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Rutherford Appleton Laboratory

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**Science and Engineering
Research Council
Central Laser Facility**

**Annual Report to the Laser
Facility Committee 1983**

RAL-83-043

Laser Division

Rutherford Appleton Laboratory

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SCIENCE AND ENGINEERING RESEARCH COUNCIL

CENTRAL LASER FACILITY

ANNUAL REPORT TO THE LASER FACILITY COMMITTEE 1983

PREFACE

The report covers the work done at, or in association with, the Central Laser Facility during the year ended 31 March 1983. It follows the same general format as last year's report except that there is an additional chapter to cover the work of the newly established Ultra Violet Radiation Facility (UVRF). There are therefore eight chapters in all, six corresponding to the six groups of the Glass Laser Scientific Programme and Scheduling Committee (GLSPSC) a chapter (chapter 2) on gas laser development and chapter 3 describing the work and development of the UVRF.

The most important single development at the CLF during the past year has been the establishment of the UVRF. The provision of more extensive facilities for University users of rare gas halide and other UV lasers has been under discussion for some time. Though originally conceived as based on the ELF laser (see previous annual reports) it soon became clear that many potential users preferred to sacrifice power for the higher pulse repetition rates available from discharge excited lasers, provided that spectrally narrow line outputs and good beam quality could be obtained. This has been done using a grazing incidence diffraction grating and double passed amplifier with a phase conjugate SBS mirror (see Chapter 3). The new facility has already proved very popular and is a rapid growth area with users from University departments of Chemistry, Physics, Biology and Engineering.

On the Vulcan glass laser the most important single development has been the installation of three additional disc amplifiers, thus completing the capital project originally approved in 1978. To do this installation

Vulcan had to be shut down for 3 months but the opportunity was taken to improve various features of the laser including the physical layout. More details are given in Chapter 1.

Chapter 2 records the successful operation of SPRITE as an oscillator both with and without injection locking. The original specification (over 200 joules in 60 ns using KrF) has been met and it is possible that this may be improved upon with further refinement. However attention has now shifted to multiplexing and preliminary work on this aspect of the final system is described in the report together with further work on UV phase conjugation and other pertinent topics.

The experimental programme using Vulcan, described in Chapters 4 to 7 inclusive, was inevitably curtailed by the shutdown of Vulcan mentioned above. However important progress has been made in various directions. Chapter 4 on laser plasma interactions describe observations of the simultaneous emission of $\omega/2$ and $3\omega/2$ harmonics of targets irradiated with green light which show that the line splitting is the same, in frequency, for the two harmonics which gives strong evidence in favour of the Raman mechanism. The same chapter describes studies of filamentation and backscatter with very much higher temporal resolution than hitherto which give additional clues regarding the nature of filamentation instabilities.

Chapter 5 covers a wide range of topics in the general field of transport ranging from continued studies of thermal smoothing to novel experiments such as the first measurements of the energy loss of α -particles in hot matter. For the latter experiment the source of α -particles was a thin walled, exploding pusher, DT filled target and the α -particle energies were determined from the tracks in a CR39 detector placed on the other side of a laser heated foil. Another set of novel observations described in this chapter concerns the study of radiative preheat by observation of the broadening of the K absorption edge: broadening dominates over a shift in the edge due to the strong coupling in the high density material achieved.

Steady progress towards understanding and quantifying various factors of importance in ablative compression is described in Chapter 6. particularly noteworthy are the studies of the growth of the Rayleigh-Taylor instability using rippled targets and the demonstration of much better agreement on compressed core sizes between theory and experiment when relatively low aspect ratio shells are driven with the higher ablation pressures available at 530 nm. These results were obtained before the laser was upgraded as described above and are therefore particularly encouraging.

The Carbon VI laser experiment remains a major activity of the Spectroscopy group whose work is described in Chapter 7. Green illumination has been used for the first time in this work, including cylindrically symmetric illumination using 4 beams. Though improved energy coupling was obtained the gain was not markedly improved and clearly problems remain. More generally X-ray spectroscopy has been used in a variety of diagnostic situations. Amongst a number of novel applications is the observation of time resolved EXAFS spectra using laser generated X-rays. This technique has considerable potential and is complimentary to the more conventional approach using synchrotron radiation.

The Theory and Computation chapter, Chapter 8, is as usual the largest. It describes work on code development and a wide range of theoretical topics, many related to the experimental programme described elsewhere. The problem of implosion symmetry, for example, is particularly topical and it is shown that the latter is more complex than appreciated hitherto. The Bangor group have made a major contribution to the theory chapter this year, particularly with respect to laser plasma interactions including the coupling of interaction processes, such as the effect of Brillouin scattering on Raman scatter. Another major topic of practical importance is radiation transport together with the related subjects of opacity and radiation preheat in laser driven implosions. Finally, in a departure from normal practice, a brief account is given of two CECAM workshops in which CLF users and staff were particularly involved.

A general feature of this year's report to which it is worth drawing attention is the growing degree of international collaboration in both theoretical and experimental work with a two way flow of staff and ideas. Not only have staff from overseas laboratories collaborated in experiments at RAL but also we report experiments carried out by UK University users in overseas laboratories.

Finally, as Laser Facility Committee members will know, this is the last annual report for which I will act as general editor. I would like to take this opportunity to thank committee members, CLF staff and University users for making my tasks as head of the Laser Division, RAL, so easy and pleasant. I have enjoyed the past 6½ years very much.

A F Gibson

CHAPTER 1 GLASS LASER FACILITY

INDEX

- 1.1 INTRODUCTION page 1.1
- 1.2 GLASS LASER OPERATIONS page 1.1
- 1.3 GLASS LASER DEVELOPMENT page 1.3
 - 1.3.1 Glass Laser Enhancement
 - 1.3.2 Pulse Generator Enhancement
 - 1.3.4 Computer Control and Data Acquisition
 - 1.3.5 Extended Clean Room Facilities
 - 1.3.6 Single Beam Disc Amplifier
 - 1.3.7 Flash Lamp Test Facility
- 1.4 TARGET AREAS page 1.19
 - 1.4.1 Existing Target Areas
 - 1.4.2 New Target Areas
- 1.5 OPTICAL DEVELOPMENT page 1.21
 - 1.5.1 Aspheric Lenses
 - 1.5.2 Uniformity of Multi-beam Illumination
- 1.6 TARGET PRODUCTION AND DEVELOPMENT page 1.24
- 1.7 DIAGNOSTIC DEVELOPMENT page 1.27
 - 1.7.1 Sub-nanosecond X-ray Framing Camera
 - 1.7.2 Sub-nanosecond Gated X-ray Image Intensifier
 - 1.7.3 Compact X-ray Streak Camera
 - 1.7.4 X-ray Pinhole Camera with Image Intensifier Recording
 - 1.7.5 Inertial Target Mounting System for Microballoons

REFERENCES page 1.37

CHAPTER EDITOR: J E Boon

CHAPTER 1 GLASS LASER FACILITY

1.1 Introduction

A number of important improvements have been made over the past year to the glass laser, its control system, diagnostics and supporting facilities such as extended clean room facilities and a lamp test facility. These developments are described in detail in section 1.3. Particularly important has been the installation of three additional disc amplifiers so that each of the 6 beams fed to TA1 now uses a separate final amplifier. At the same time a second beam has been provided for TA2 (section 1.4) which will considerably increase scheduling flexibility. Looking further ahead work has begun on the development of a larger aperture disc amplifier (section 1.3.6), new target areas are under construction (section 1.4.2) and lens development is continuing (1.5.1).

Diagnostic instrumentation development is traditionally done in Universities and this has continued (sections 1.7.1 to 1.7.4 inclusive). Section 1.7.5 describes a joint development between RAL and MPQO, Garching, which may prove a very important contribution to the problem of target illumination uniformity, which is discussed in section 1.5.2.

A F Gibson (RAL)

1.2 Glass Laser Operations

The Glass Laser Operations Group has had a busy year starting with the first operational period with the new synchronised long and short pulse oscillators and ending with a major shutdown to enhance the laser system (section 1.3.1). Of the 52 weeks in the year, 33 were available for target area shots, 6 for laser maintenance and 13 for the laser enhancement shutdown.

Considerable improvements have been made to the laser availability during

the experimental periods with the laser now running from 08.30 to 17.00, 4½ days a week. This has mainly been achieved by improvements to the alignment stability such that adjustments are unnecessary over periods of weeks.

The two major causes of downtime have been the oscillators during the early part of the year, and the beam apodisers. Oscillator reliability has improved dramatically with operating experience and should improve further with the introduction of developments covered in 1.3.2. The apodiser problem centred on water evaporation from the glycerol/water/copper sulphate mix raising the operating temperature gradually to the point where copper deposition on the windows occurred. This has been successfully overcome by using mixes with lower initial operating temperatures giving a slower initial rate of temperature increase, and regular changes to the fluid before its index match temperature becomes too high.

Other work during the year has included the installation of a new laser control system (section 1.3.4) and the synchronisation of the laser to a Z-pinch discharge.

The year ended with a shutdown to enhance the laser capability to 1.2kJ in six beams, and make a number of changes to the laser staging based on operating experience. Initial operating experience with the enhanced system has proved very encouraging in terms of reliability and operational simplicity. Outputs per beam of 0.6 TW in 70 ps and 150 J in 1 ns have so far been delivered to target areas with long pulse energy being limited by target area requirements only. Full information on the enhancement is in section 1.3.1.

N Allen, D Gleed, A Raven, P Rumsby and R Wyatt (RAL)

1.3 Glass Laser Development

1.3.1(a) Glass Laser Enhancement

Enhancement of the glass laser output by increasing the number of output disc amplifiers from 3 to 6 in the six beam system represents the final phase of the enhancement programme started in 1981. Operating experience with that system however dictated the need for more extensive changes to the system to improve the reliability and output capability. Specific problems were

- (i) The double pass output scheme, while economical in terms of amplifiers per joule of output, needed an output polariser in the highest fluence part of the beam. Being one of the weakest components for damage threshold, the output was being limited by polariser damage. Although single polarisers had been run at 220 J output, the financial cost of damaging 6 output polarisers in a single shot resulted in very conservative operating levels of ~ 120 J/amplifier.
- (ii) Introduction of the synchronised oscillator system has led to the need to alternate short and long pulses down the same rod amplifier line and the consequent need to change amplifier gains between shots. This lead to the risk of firing short pulse shots at the gain settings for long pulses.

The first of these problems was overcome by changing all output amplifiers to single pass and using quarter wave plates coupled with the polarising properties of the disc amplifiers and output mirrors to provide backwards isolation. This requires an increase in drive into the discs by a factor ~ 10 (single pass disc gain) and a consequent redesign of the driver stages. The system "fuse" is now the output polariser of the double passed driver disc amplifier which sees approximately the same fluence as at the output of the six discs. However the damage risk is now reduced to one polariser only. The second problem was solved by dedicating a rod amplifier chain to a particular pulse length - long or short - with pulse switching into the output amplifiers occurring at the end of the rod

chains.

1.3.1(b) System Design

The system layout is shown in Figure 1.01. The oscillator system remains unchanged and consists of long and short pulse oscillators synchronised together. Both switched out pulses are amplified by two 9 mm rod amplifiers, apertured and spatially filtered before being injected as expanding gaussian beams into the two rod amplifier chains. After further amplification the beam is apodised (AP) in a profiled and index matched absorbing cell and then continues expanding through amplifiers of increasing aperture. Computer controlled wave plates (WP) (section 1.3.4) in front of the Faraday rotators (FR32) allow the chain energies to be set. Waveplates are also used to run circular polarisation through the rod amplifiers. This ensures that any birefringent loss is axially symmetric and removes the "squaring up" of the beam shape at high repetition rates. ASE from the rod amplifiers and prepulse are suppressed by two 50 mm aperture pockel cells (PC). These are driven with a 40 ns pulse from a Cardon Instruments thyatron pulser to provide a 25 ns flat top open gate. This also serves to provide backward isolation with the gate closing before the first back reflection returns (from leakage from the double pass discs). The pulser is electrically triggered from the oscillator Q-switch signal and because of the inherent stability of the build up time of the synchronised oscillators (< 1 ns jitter) the Pockel cell gate is essentially jitter free with respect to the pulse.

Fifty per cent splitters after the two Pockel cells split each beam into two to provide long and short input pulses into both disc amplifier chains the selection being made by two movable mirrors. 76 mm amplifiers compensate for the split and the beam is then expanded to a 105mm parallel beam before entering the disc amplifier chain. Currently the driver disc is fed by a beam expander, but this is shortly to be replaced by a vacuum spatial filter. Further waveplates in front of the driver and single beam discs allow variations of the output of the two arms. Because both these discs now run at comparable outputs of 200 J/0.5 TW the opportunity has been taken to provide a second beam to the single beam target area from

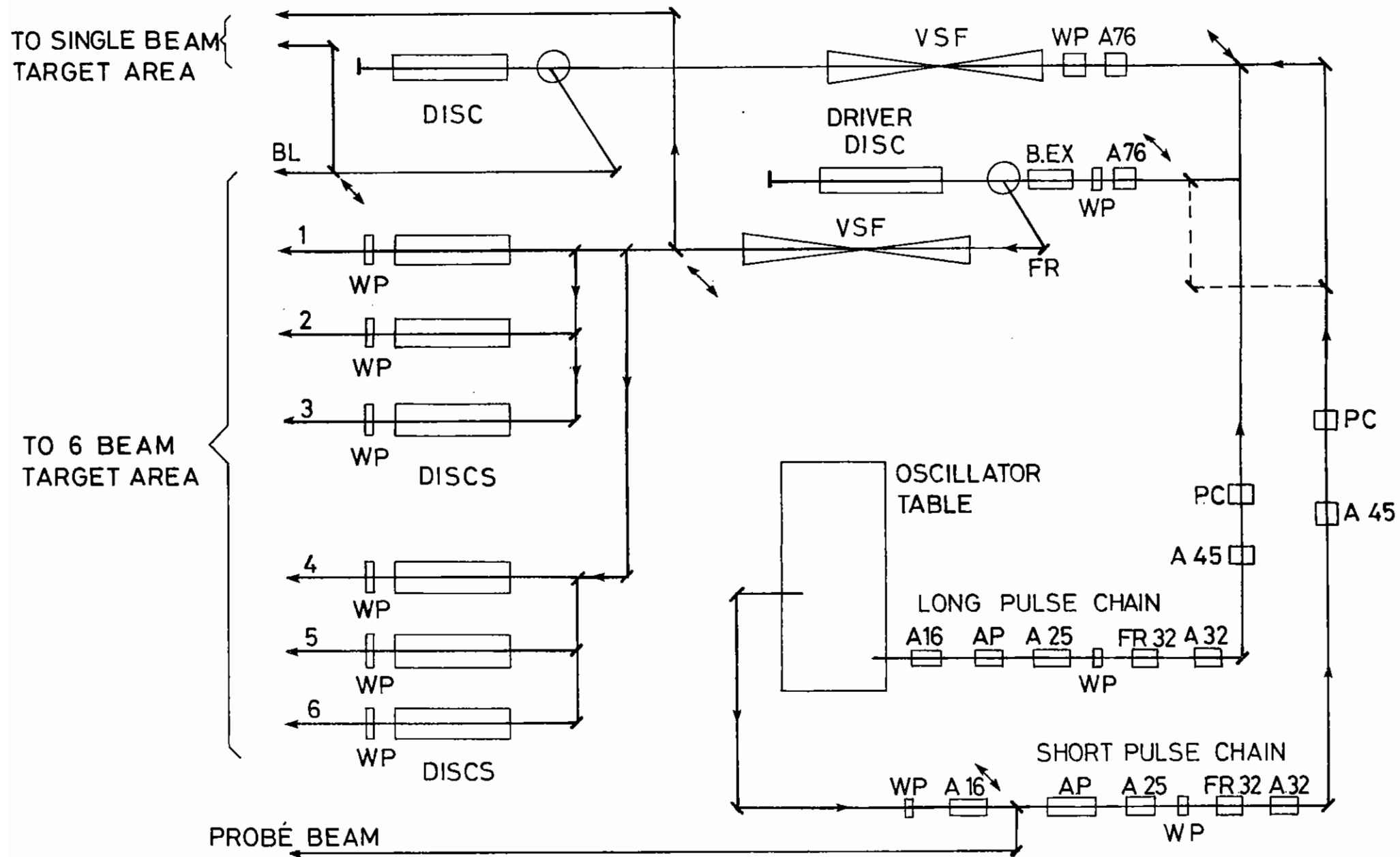


Fig. 1.01 Glass laser system after recent enhancement.

the driver disc, allowing backlighting experiments to be carried out and so easing the load for this type of experiment on the six beam target area. When not diverted to the single beam area the driver disc output is divided into six and time delayed for synchronism at the target before being further amplified by the output discs. Quarter wave plates at the output provide retro isolation for infrared experiments but can be removed for green operation.

The single beam disc is double passed as before but with the polariser now operating in reflecting mode output. Damage measurements (1.01) on the polarisers give 6 J/cm^2 damage level for p and $\sim 20 \text{ J/cm}^2$ for s-polarisation indicating that s-polarised output is to be preferred. The output from this disc can be switched to provide either the main beam for the single beam target area (TA2) or a backlighting beam for the six beam area (TA1).

1.3.1(c) Commissioning

In routine use the early waveplates are set to provide the maximum allowable output from the system (200 J/0.5 TW per beam) with the later waveplates fully open. These later waveplates are then used to change the output from 100% to 0% of maximum. No problems have been encountered with these items or with the newly introduced pockel cells. Outputs of 3.2 TW in 70 ps have been routinely achieved in the six beam system (no long pulse experiment has taken place since the upgrade) and 150 J in 1 ns in the single beam limited by experimental requirements only. It is hoped to take this up to 200 J during the coming experimental period. Beam balance can be held to the $\pm 10\%$ level for infrared operation and comes mainly from gain variations in the output discs. This arises from a combination of newer formulation glasses having higher gains at fixed doping plus the effect of repeated repolishes gradually reducing the disc thickness by up to 2 mm in 30 mm for some of the oldest discs in use. We are currently examining the possibility of polarisation tuning the split ratio to the

six beams to compensate for the amplifier gains and bring the balance with 5% total excursion.

N Allen, D Craddock, R Day, M Forster, D Gleed, P Gottfeldt, W Lester, A Raven, C Reason, P Rumsby, S Sunderland, R Wellstood and R Wyatt (RAL)

1.3.2 Pulse Generator Development

This year has seen three main areas of activity - finalising the previously reported work on the synchronised laser oscillators; developing the "intelligent attenuator", a novel device for improving long and short term pulse energy reproducibility from pulsed lasers; and developing a device to reliably maintain single longitudinal mode performance from the long pulse oscillators, by feedback control of the optical resonator length.

1.3.2(a) Synchronized Oscillators

The synchronized oscillator system (1.02) has now completed more than a year of operation, and has proved reliable and effective. Only two significant long term reliability problems emerged from the extended period of running which it has now seen. One was degradation of the circulating coolant in the short pulse oscillator head, which under the $\sim$ nine millisecond duration flashlamp pulses proved prone to algae infestation. After experimenting with commercial biocides and various coolant liquids, a 50% mixture of purified water and ethanol was found to at least contain the problem at a level where the system now runs for several months without noticeable coolant degradation. The use of cerium doped flashlamps was also found to help considerably, presumably by reducing the amount of ultraviolet radiation reaching the coolant.

The other problem was an unacceptably short flashlamp lifetime ($< 10^3$ shots) in the short pulse oscillator, under the loading of the 9 ms, 120 amp driving pulse required for actively mode locked and actively Q-switched operation. The lamps did not fail explosively, but started to give out an audible "click" associated with a ruinous breakup of the

prelase and a serious degradation in the amplitude and timing of the Q-switched output train. By replacing the original pulsed type Xenon flashlamps with CW-type krypton flashlamps at higher pressures the lifetime has been increased to $> 10^5$ shots. This is acceptable, with our current low duty cycle on this oscillator (it is not run continuously between shots) but efforts are continuing, in partnership with the flashlamp manufacturers, to improve further.

1.3.2(b) The Intelligent Attenuator

The "intelligent attenuator" (1.03) or self setting attenuator is an active electro-optic filter which we have developed to dramatically improve the pulse to pulse energy reproducibility of pulsed laser systems. It is shown schematically in Figure 1.02. It was originally conceived as a linear device, to reduce the $\sim \pm 10\%$ typical long term drift of the synchronized oscillator systems output energies down to $< 1\%$, for precise predictability of laser output energy. During development it became clear that a non linear mode of operation existed, capable of pulse energy stabilization over a very wide range of input energies, in an optical analogue of the electrical Zener diode. This makes the device rather versatile for more general applications.

(i) Method of Operation

The input laser pulse energy is sampled by a photodetector, which feeds through the combination of a pulse forming network (PFN) and a high voltage amplifier (HVA) to apply to the Pockels cell a voltage pulse which is related to the input pulse's energy. By a correct choice of parameters the assembly can act as a fast real time feed forward energy stabilizer, by applying a higher attenuation to a higher than normal energy input pulse, and vice versa, giving output pulses of quasi-constant energy. The PFN essentially stretches the photodetector pulse in time to the extent that it comfortably brackets the input laser pulse, and provides a constant attenuation throughout the pulse duration.

The high voltage amplifier must provide kilovolts onto the Pockels cell to

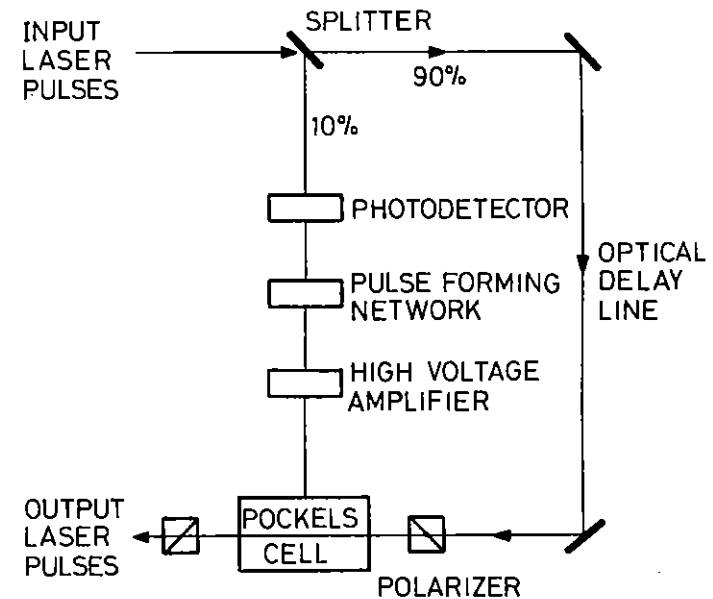


Fig. 1.02 Schematic of the intelligent attenuator.

effectively modulate the 1.05 μm laser pulse. This, in conjunction with the < 10 ns response time required for sensible processing of the same laser pulse, lead us to consider first the case of a simple linear amplifier, rather than more complex digitally programmed circuitry which could provide excellent response characteristics but would be too slow. For the linear amplifier case the transmission of the device can be represented by (1.03):

$$E_{\text{out}} = E_{\text{in}} \cdot \cos^2 \left(\frac{D + K \cdot E_{\text{in}}}{V_{12}} \right) \quad (1)$$

where E_{out} and E_{in} are the output and input pulse energies, K is a constant, D is a DC bias applied to the Pockels cell (to reduce the output swing required from the high voltage amplifier) and V_{12} is the half wave voltage of the Pockels cell.

The degree of output pulse energy stabilization given by this simple relationship is shown in Figures 1.03 and 1.04. It is rather good for fairly small fluctuations about the operating point (the peak). For instance input energy fluctuations of $\pm 10\%$ are reduced to $\approx (+0, -1\%)$ at the output. Also the peak transmission is seen to be quite resonable, up to 52%. The use of a DC bias, to reduce the voltage swing demanded from the linear amplifier, is seen to also give slightly better stabilization properties, but at the expense of a somewhat reduced transmission. "Optical Zener diode" mode operation is obtained by arranging that the high voltage amplifier begins to saturate beyond the peak value of E_{out} , so that the rate of fall of E_{out} vs E_{in} beyond the peak is reduced, or even eliminated.

(ii) Experimental Demonstration

The demonstration of the device used the equipment shown in Figure 1.05. In practice of course the difficulty is the combination of a photodetector, PFN and amplifier which are capable of kilovolt responses with nanosecond risetimes. The connection of a photoconductive silicon switch, run at low optical input intensities in the linear part of its

characteristic, and a single stage planar triode amplifier proved very adequate. Voltage outputs of ≈ 1.2 kV (linear) and 2 kV (at 20% saturation) were achieved with a risetime of < 5 ns. Also a simple combination of a charge time and a shortened stub was all that was required for the PFN, because of the properties of the silicon switch.

Figures 1.06 and 1.07 show the experimental results verifying intelligent attenuator action under different parametric conditions (DC bias, silicon switch filter, input energy) varying from close to linear, in cases (a) to optical zener diode mode in cases (d). Notice how this latter mode of operation gives the output pulse energy stabilized with a band whose extrema are $+5\%$ and -8% while the input energy varies by more than a factor of six!

The intelligent attenuator concept is very versatile, and can be applied over a broad range of wavelengths, pulse duration and energies. We plan to install the device on each beamline of the laser in the near future.

1.3.2(c) Single Longitudinal Mode Resonator Stabilization

During the year we began a project to feedback stabilize the optical resonator length of our single longitudinal mode (SLM) oscillators, including the long pulse generator of the synchronized oscillator system. This would eliminate the occasional drift into multi mode operation (rippled pulse temporal profile) which in the absence of feedback stabilization is inevitably suffered by all SLM oscillators due to thermal and mechanical drift of the resonator's optical path length and the mode selector's resonance. Two elegant techniques for accomplishing this feedback control have already been described (1.04)(1.05) but were not applicable to our case. Hanna (1.04) has used the "thermal self sweeping" inherent in a Nd:YAG oscillator to avoid multimode failures, but the technique is difficult to apply to our Nd:YLF oscillators (matched to the phosphate glass of the amplifier chains) because of the much lower self sweeping in this athermal material. Kuizenga (1.05) has developed a prelude sampling technique, but again the implementation is very much more difficult with YLF and the technique is not applicable to "short prelude"

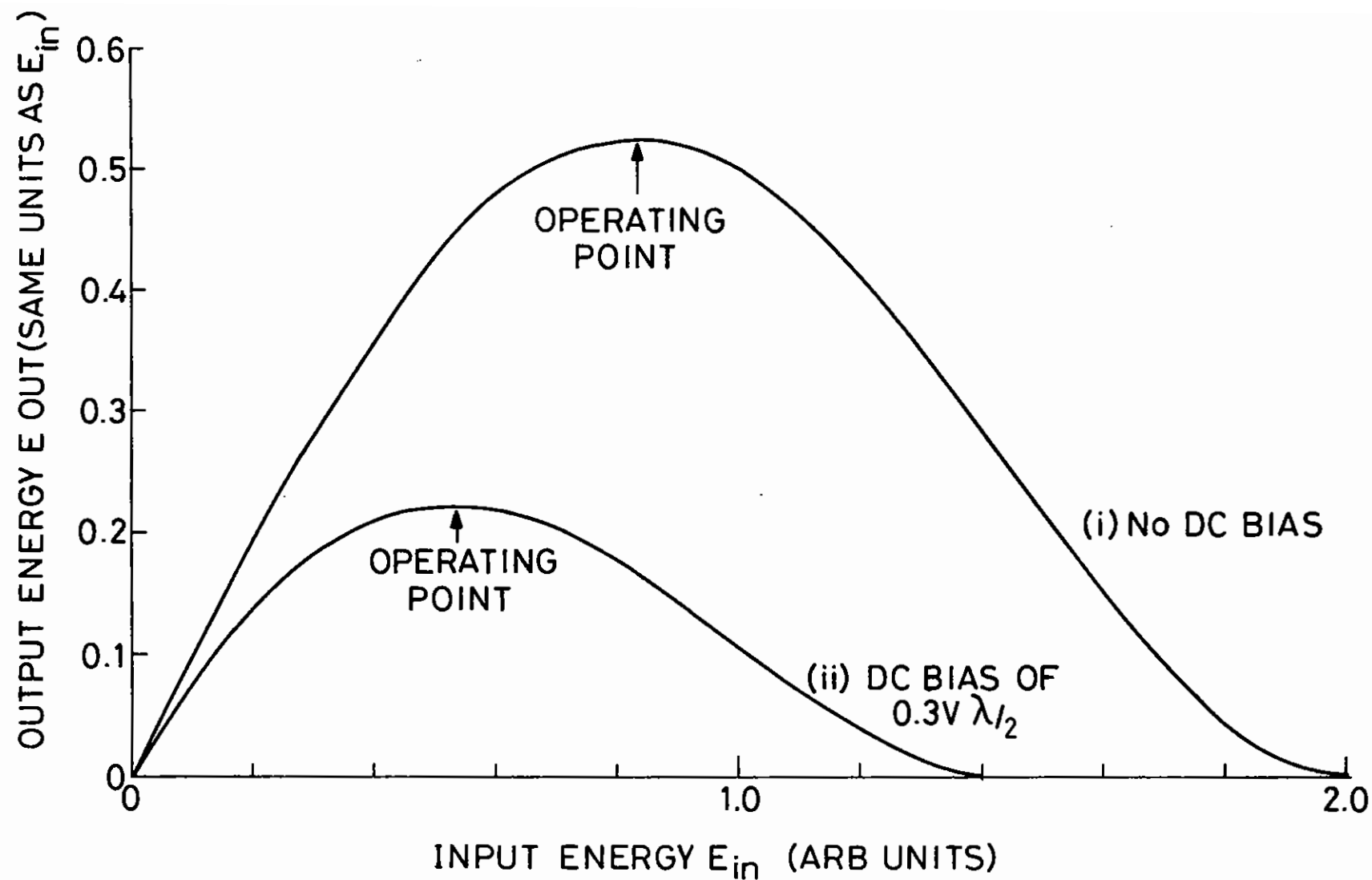


FIG 1.03 THEORETICAL INTELLIGENT ATTENUATOR CHARACTERISTICS FOR CASE OF IDEAL LINEAR AMPLIFIER

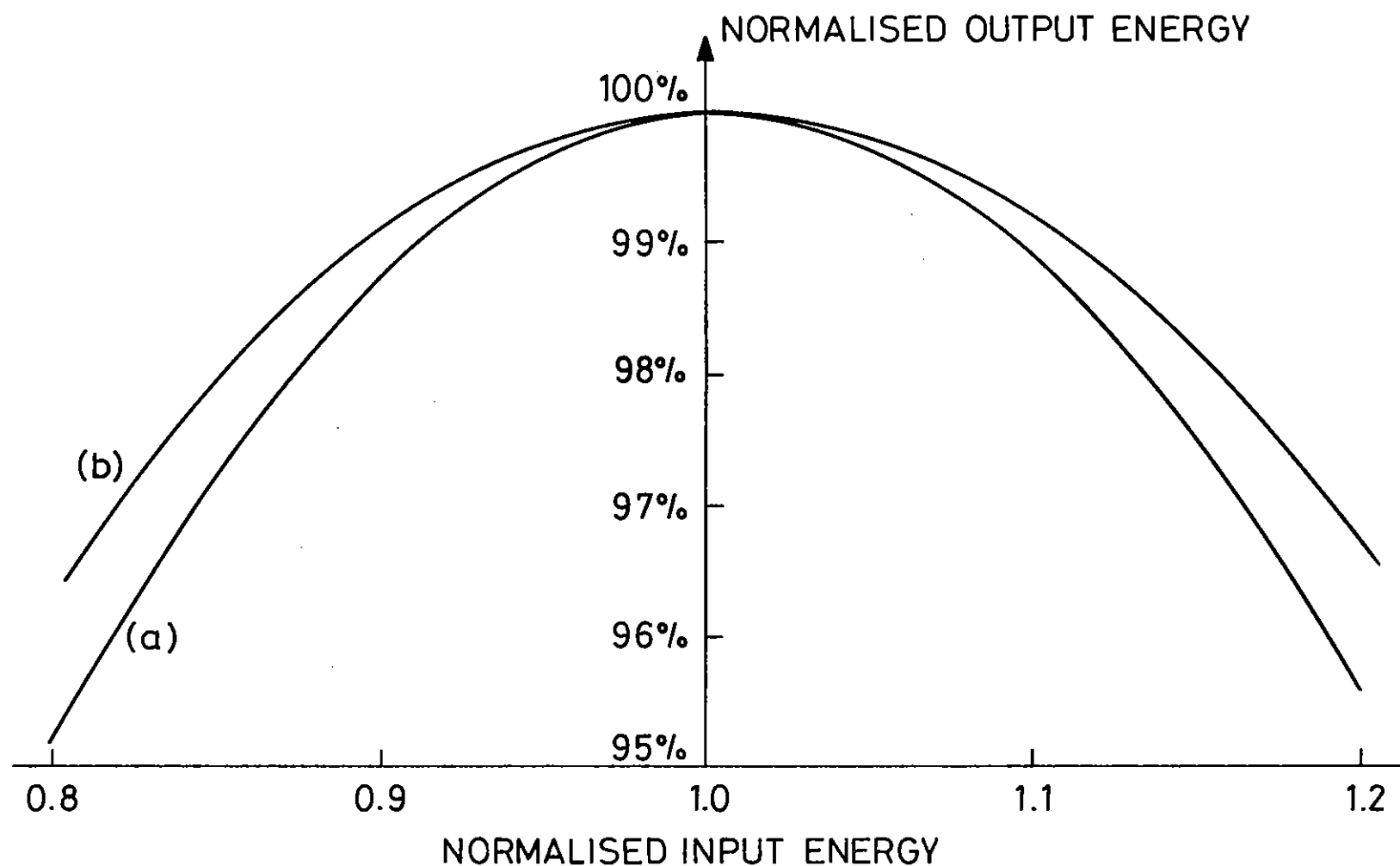


FIG 1.04 THEORETICAL INTELLIGENT ATTENUATOR CHARACTERISTICS
EXPANSION OF FIGURE (2) CURVES AROUND THE
OPERATING POINT IN NORMALISED UNITS

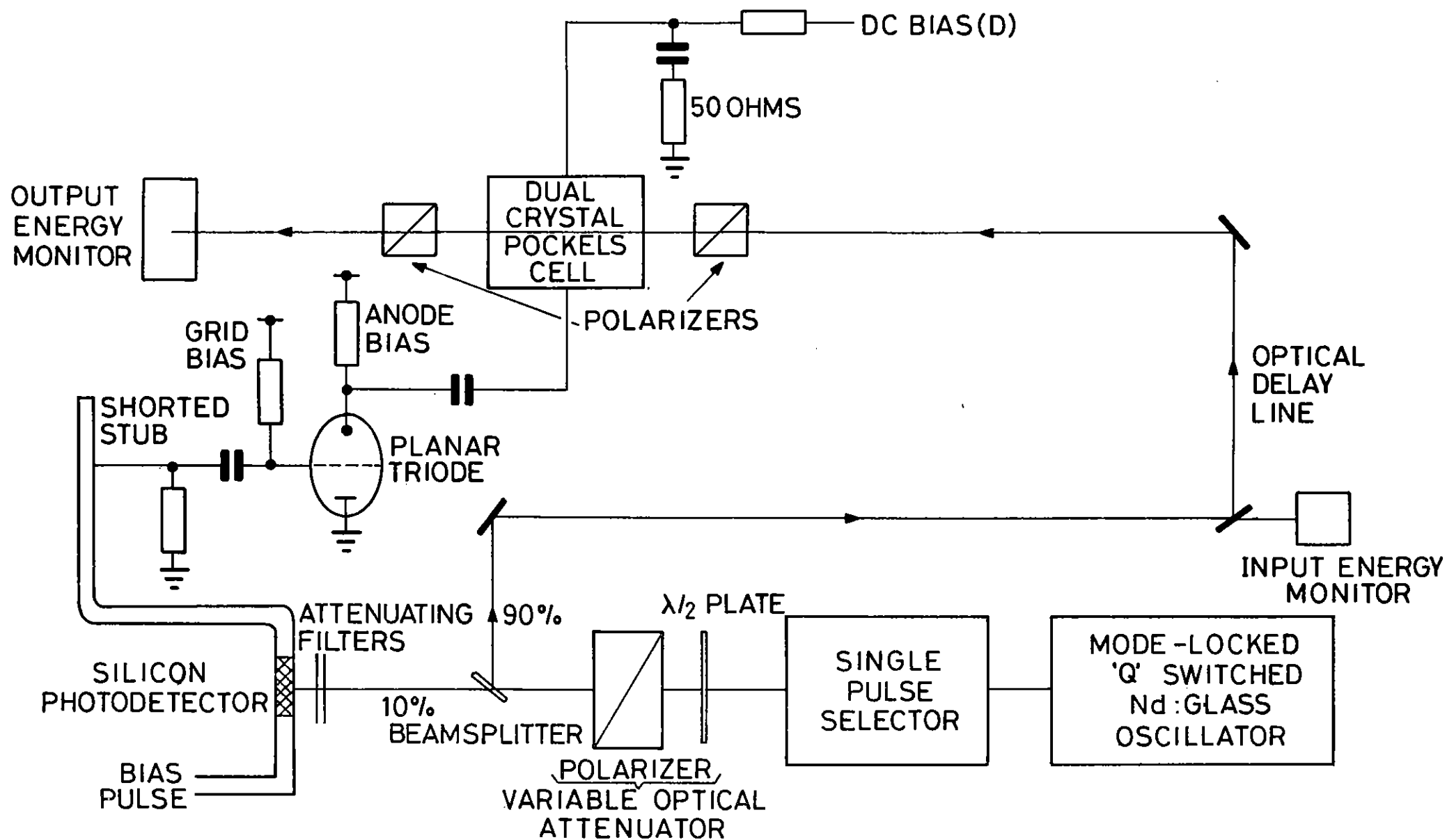


FIG 1-05 EXPERIMENTAL ARRANGEMENT

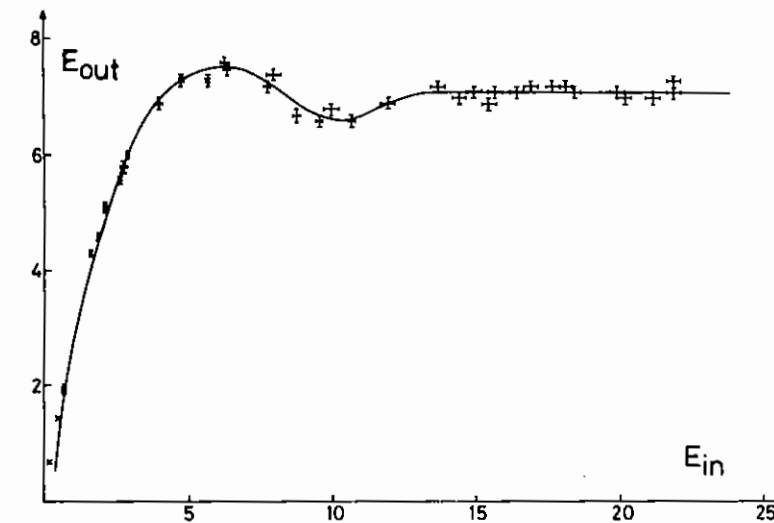
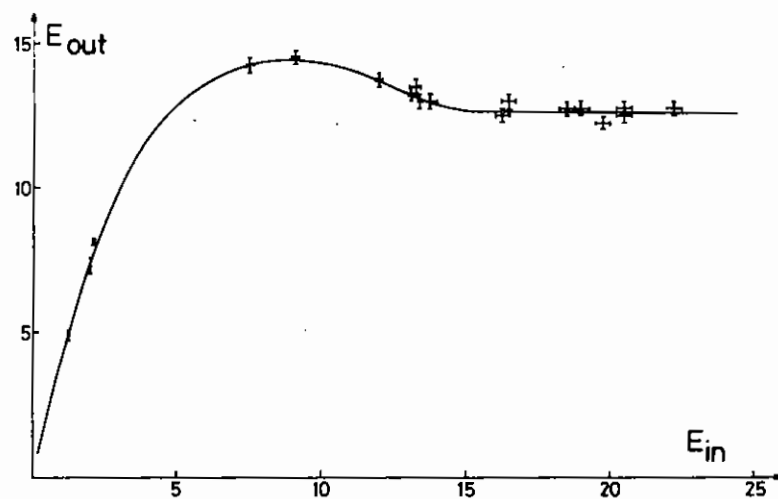
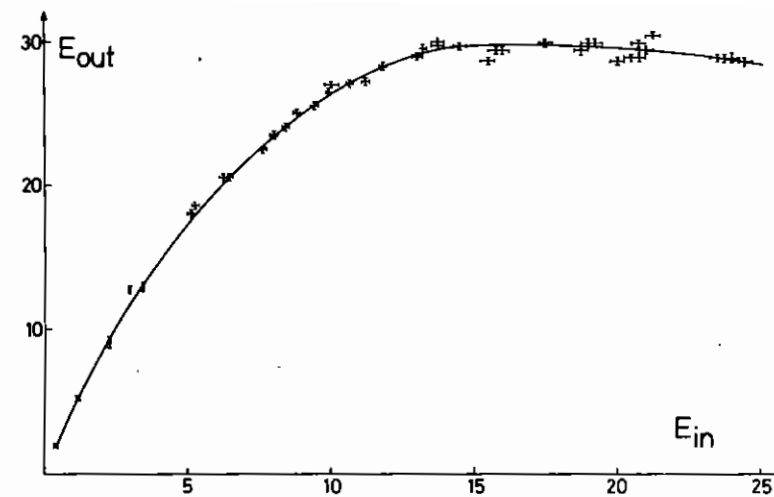
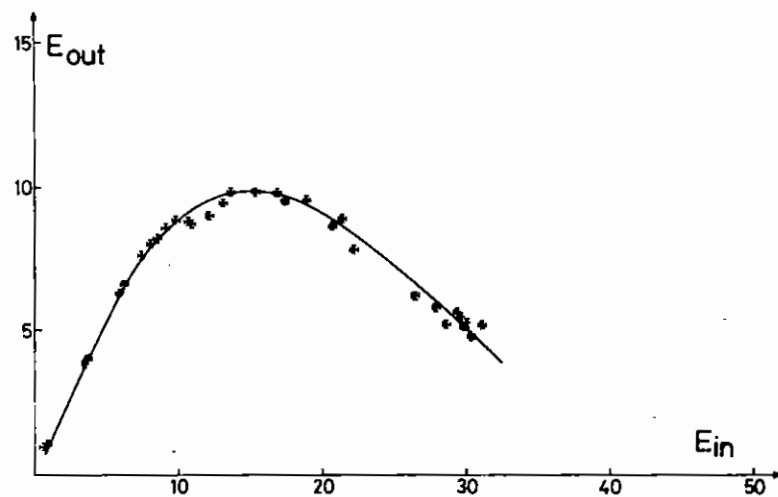


Fig. 1.06 Experimental intelligent attenuator characteristics.

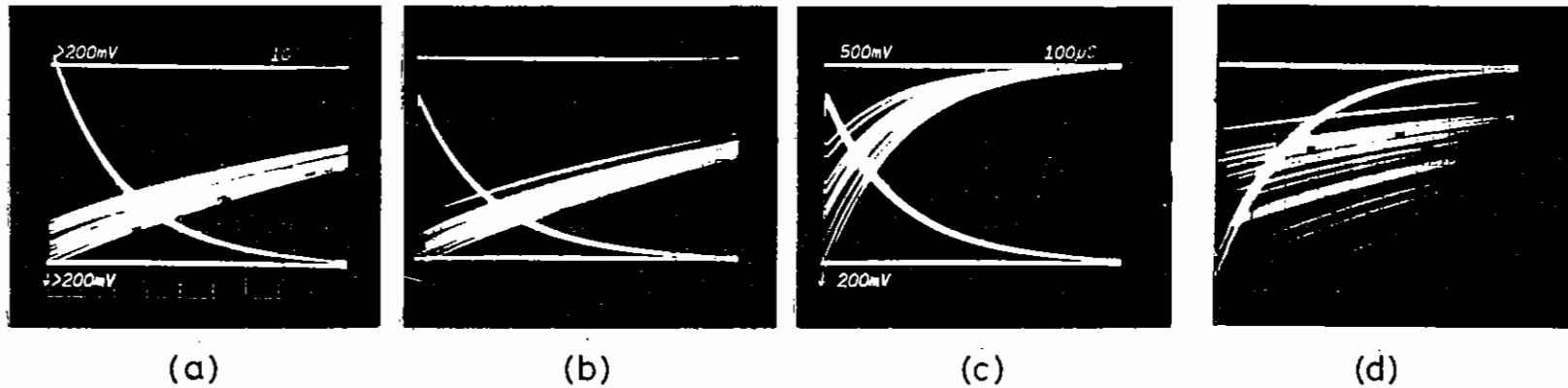


FIG 1-07 MULTI SHOT OVERLAY OF INPUT AND OUTPUT ENERGY MONITOR PHOTODIODE TRACES, ILLUSTRATING THE INTELLIGENT ATTENUATOR ACTION ($\Delta E_{\text{out}} \ll \Delta E_{\text{in}}$)

SLM oscillators such as are used here.

We propose to use a technique of sampling either the pulse envelope or pre-lase signal for mode beating using electrical filtering, followed by a step adjustment of the piezo-driven back mirror when mode-beating is detected. Using such a sample and adjust technique the oscillator should be able to self adjust itself to single mode operation during the charging period to a shot (the oscillator cavity is mechanically blocked outside this period for safety reasons).

Tests with sampling the filtered output of the pulse envelope have demonstrated the feasibility of such a system, but problems are encountered with attempting to detect low modulation levels because of the presence of high frequency components of the 25 ns pulse envelope itself. We are currently looking at detection from the much longer period pre-lase signal as used by Hanna (1.04). The possibility also exists of replicating the techniques of reference 1.04 by using the piezo stack to mechanically sweep the cavity during the pre-lase period.

A Brett, A Raven, M White and R Wyatt (RAL)

1.3.3 Laser Beam Diagnostics

The diagnostic system which is under development is intended to be modular in form so that measurements can be made using standard equipment. For monitoring the six beam outputs there are three tables on which the diagnostic equipment is attached. The other monitoring points use the same equipment but the configuration is different due to space constraints. At present in all the positions where the beam is being monitored, a fraction of the beam is sampled by using a lossy mirror. However when full monitoring is installed this will not necessarily be the case, for example at the polarisers in the system.

Three main measurements will be made to determine the performance of the laser system, these are beam energy and pulse length at various points in the system together with some representation of the intensity profile of

the beam. Each of these will now be discussed in more detail.

1.3.3(a) Energy monitoring

Ten energy monitors are now installed and working in the laser system. They have given consistent results with accuracy of the order of one per cent. As well as monitoring each of the six beam output and the backlighting disc output, energy measurements are made at the output of the driver disc and also near the end of the two rod amplifier chains.

The hardware comprises two distinct parts; the monitors themselves and the data acquisition system which resides in a CAMAC crate. A description of the Data Acquisition System is given in 1.3.4(c).

The energy monitors utilise a volume-absorbing disc calorimeter (shown in Figure 1.08) to detect a pulse of laser energy, by absorbing 95% of the pulse.

The calorimeter can be calibrated by applying measured amounts of electrical energy to the heater coil wound around the glass disc. Using this method, the sensitivity of the calorimeters was calculated to be of the order of 160 mJ/mV. When in place at the output of the laser system, one of these calorimeters will receive a fraction of the total laser energy, amounting to a maximum of a few Joules, thus giving an output voltage of the order of a few tens of millivolts. This size of signal is too small to drive the Data Acquisition System, which requires a signal magnitude of the order of 1 volt. A pre-amplifier (with gain variable between 175 and 440) has therefore been included in the energy monitor to produce such a signal. A further stage of pre-amplification (gain 10), is necessary in the monitors situated at the output of the short and long pulse rod amplifier chains, to attain the 1 volt signal level.

The actual construction details of the monitors are shown in Figure 1.09. The outer casing is an anodized metal tube, which slides over the inner assembly, consisting of two circular metal discs joined together by two metal bars. The front disc has a glass window in it, through which the

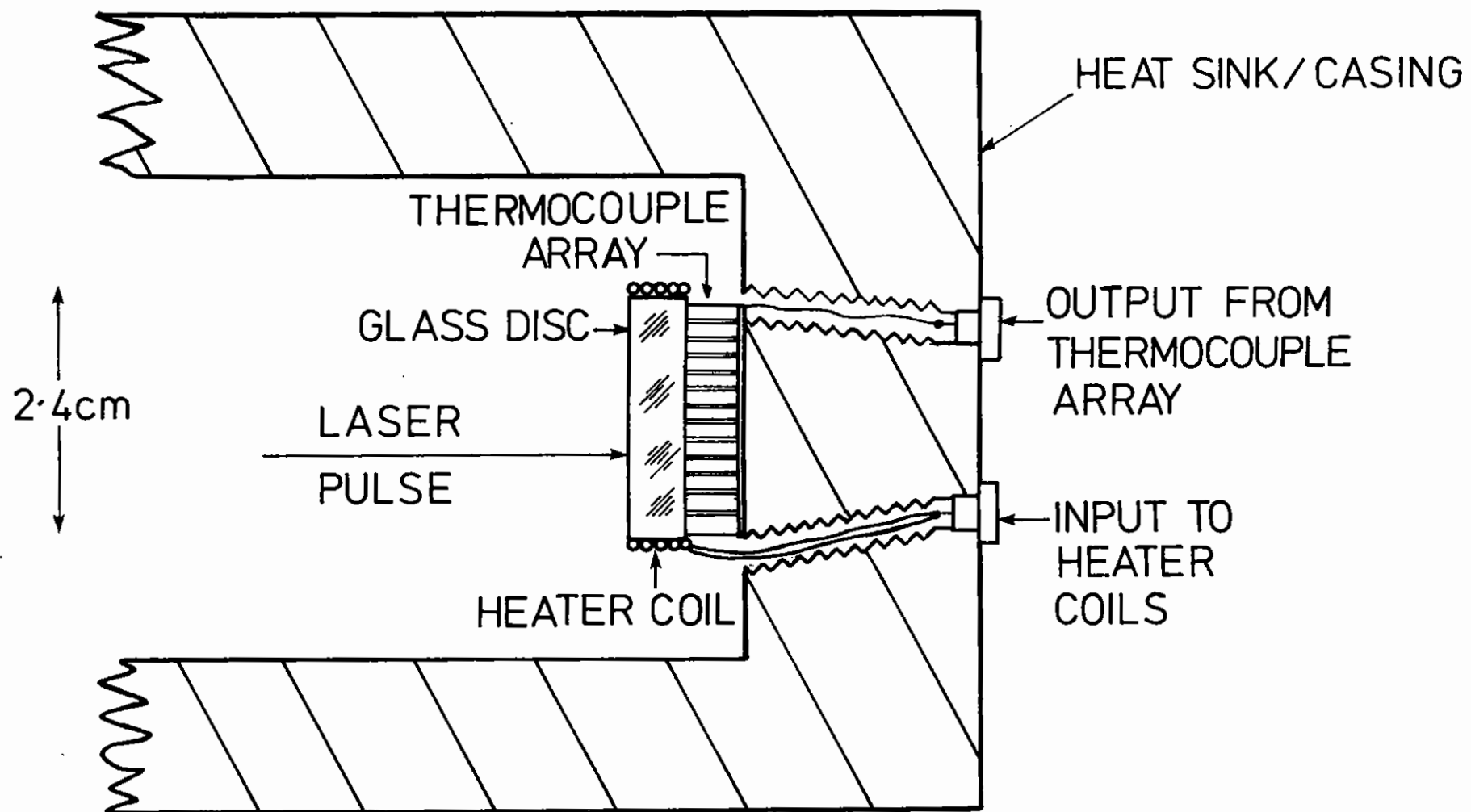


Fig. 1.08 Detail of the volume-absorbing disc calorimeter.

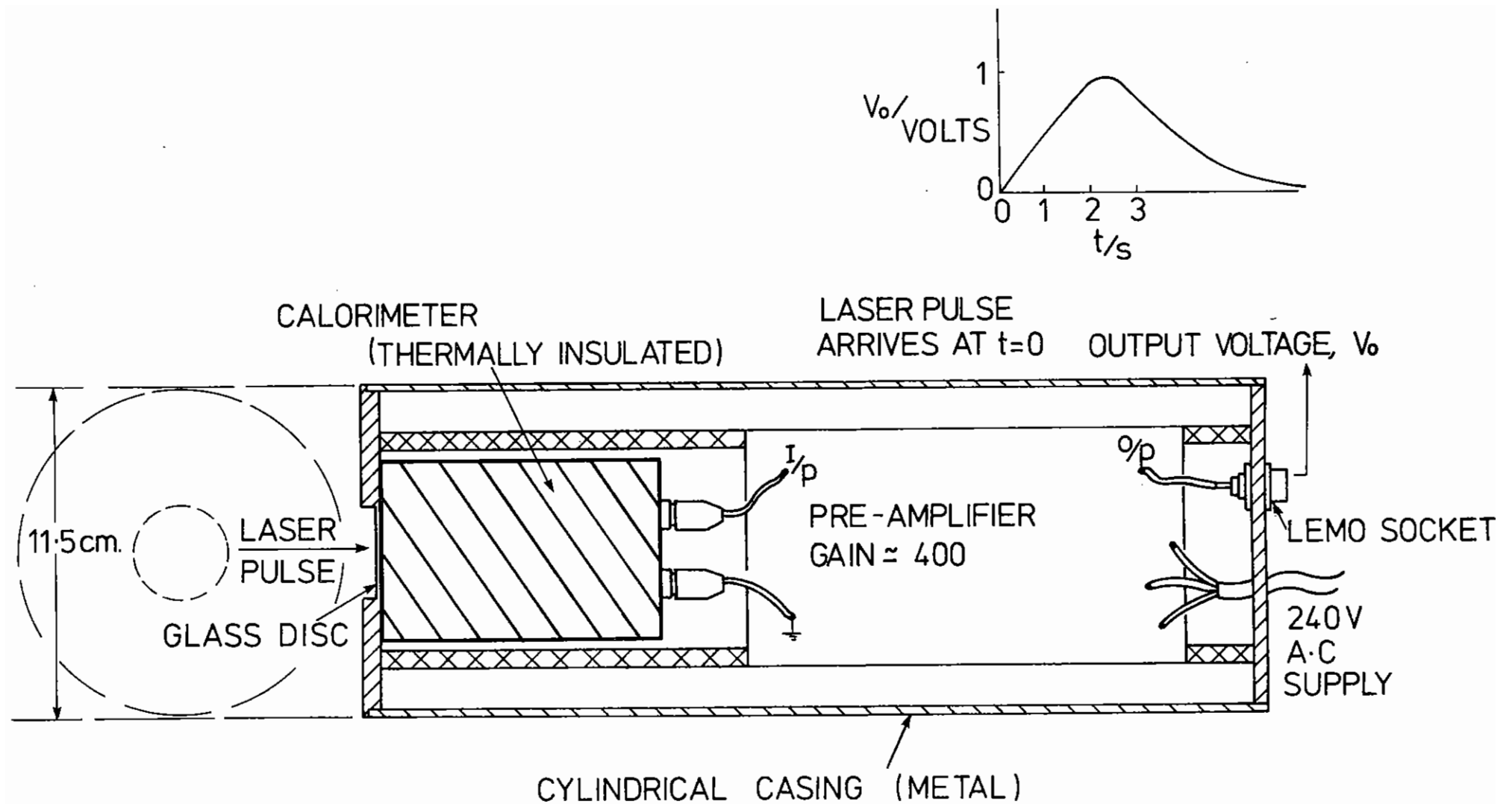


FIG 1-09
DIAGRAM OF THE CALORIMETER BASED ENERGY MONITOR

laser pulse passes, onto the calorimeter (which is itself thermally insulated using "bubble pack", to minimise variations in output due to changes in ambient temperature). The pre-amplifier is mounted on the bars, and its output is connected to a lemo socket on the rear disc. The monitor is connected to the Data Acquisition System from this via a twisted-pair cable. Each monitor pre-amplifier has its own individual mains operated power supply.

The calorimeter heater coil is also connected to two sockets on the rear disc (which are not shown on Figure 1.09), so that the monitor may be tested independently of any laser input. For this purpose, a battery powered circuit has been built, which provides a pulse of known electrical energy; enabling any monitor to be rapidly and easily tested.

1.3.3(b) Pulse Length Measurement

A Hadland streak camera has been ordered to be used in the laser room for measuring the pulse lengths of the beams. The camera will be situated near the six beam outputs and fibre optic cable (BICC type 701423002) will be used to relay to it beams which have been sampled at various points in the system. One test has been performed which involved sending the short (70 ps) pulse oscillator into TAl via 70 m of fibre optic cable to measure the pulse length on the target area streak camera. No noticeable degradation of the pulses occurred. When used in the laser room much less cable will be required.

1.3.3(c) Beam Profile Imaging

It is intended to record images of the beam profile from each of the outputs. This will involve the use of solid-state cameras and frame stores which digitise and store a single television frame. The picture information will then be archived onto magnetic tape. This system will also be used to store streak images showing the temporal profile of the beam.

Cameras from two manufacturers were tested using the Nd:YAG laser and also

oscillator pulses. As shown in Figure 1.10 the gamma of some of the cameras changed when pulsed light was used. The General Electric TN2200 camera was a low resolution model (128 x 128 PIXELS) and so could not be used even if the gamma of 1.25 was acceptable. Therefore the GE TN2506 appears to be the best choice for a camera. If further tests show that it is satisfactory it will be the standard camera for use with diagnostics.

Several manufacturers frame stores have been investigated for possible use. However because of the nature of the laser an expandable system is required which will enable many cameras to be used. As a result of this constraint together with financial ones a frame store which is marketed by Hakuto has been chosen. This device stores a single frame in its memory when given a command and then can transfer the digital data to the GEC 4080 for archiving or analysis. Many memory channels can be used and these are managed by an LSI-11/02 microcomputer which organises the transfer of the data to the main computer.

A Brasher, D Gleed and A Raven (RAL)

1.3.4 Computer Control and Data Acquisition

1.3.4(a) General Development

The GEC 4080 computer system has been further enhanced during the year to extend existing facilities and provide new control and data acquisition functions.

As a result, response to background tasks, such as controlling the micro-densitometer and in particular programme development has become ever slower. This is largely due to the time the computer spends swapping programs and data between central processor memory and disc. The response has therefore been improved by increasing processor memory from 128 Kbytes to the maximum for the 4080, of 256 Kbytes. The whole store has now been implemented in semiconductor technology. The increased cost, over simply adding another 128 Kbytes of core, has been offset against significant savings in annual maintenance charges.

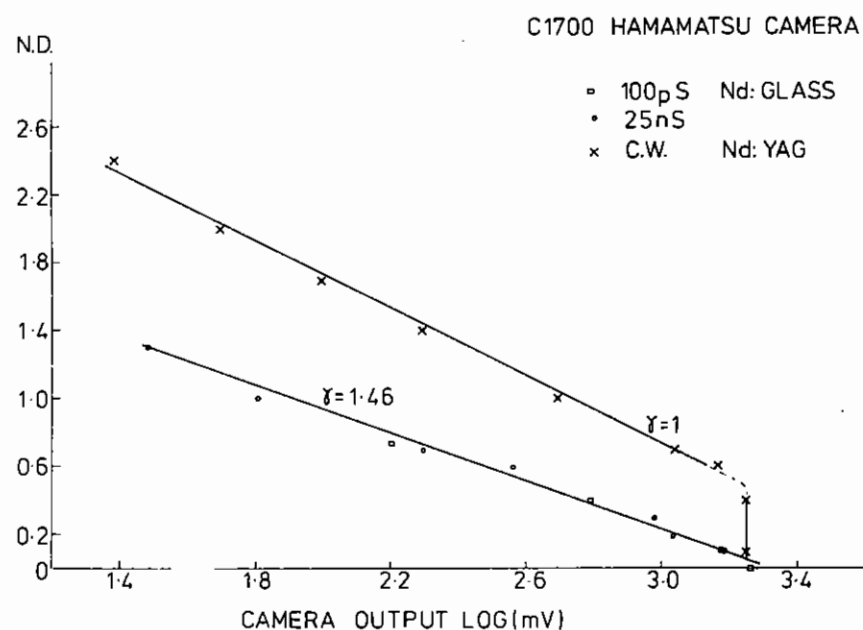
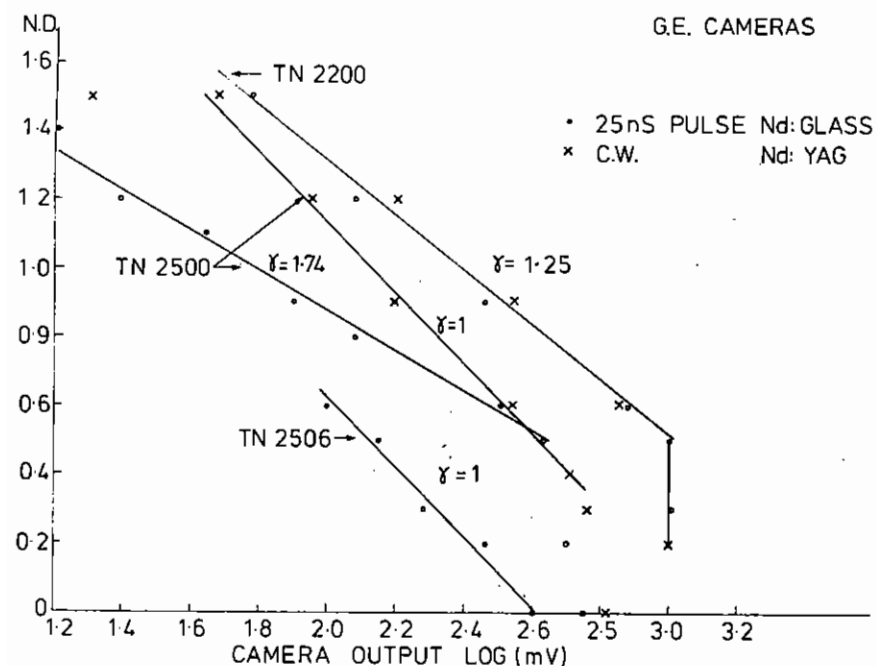


Fig. 1.10 Results of camera tests.

The main laser control program again required updating to accommodate the recent laser upgrade. In an effort to avoid gross misoperation, facilities have been included for checking the laser configuration against a "standard" set of configurations and the laser cannot be fired unless it complies. To record or modify the standards requires a password not available to the operator.

Stand alone programs have been developed during the year to analyse EXAFS data. The microdensitometer program has been extended to cope with large area scans in connection with gas laser beam profile analysis.

1.3.4(b) Control Console

A new console has been installed to provide a centralized laser control system. It houses a colour TV monitor, computer terminal, waveform digitizer monitors, a touch screen and communications systems. The colour TV monitor normally displays the laser layout but is used to show capacitor voltages during the charging sequence. It is also available for other computer generated displays. The touch screen is made of glass and is mounted in front of a TV monitor. It provides a 4 x 4 matrix of touch-sensitive areas acting as push buttons. The function of each button and the message displayed behind it are under computer control. Using this combination in a hierarchical manner a comprehensive set of laser controls will be provided. Latching circuits have been installed to allow computer controlled devices to be switched by pulses from the computer; each pulse simply changing the existing status. Manual control of the switch has also been provided for test and maintenance purposes. Using this technique, together with a status indication to the computer means that computer/manual switches are not required and the laser status quo can be maintained over computer power up/down and at initial programme loading time.

It is proposed to add other monitors to the control console to display the laser gain profile and beam spatial energy distributions.

1.3.4(c) Energy Monitoring

All energy monitoring on the glass laser now uses calorimeters (see 1.3.3(a)); replacing the light diodes used previously. A CAMAC unit has been built to read the calorimeters into the computer.

The signal is amplified at the calorimeter to about 1-volt and transmitted via shielded twisted pair cable to the CAMAC unit which can accept up to 12 inputs. The unit is micro-processor controlled and incorporates a 12 bit hybrid data acquisition sub-system comprising a multiplexer, instrumentation amplifier, sample and hold circuit and 12 bit ADC in one package. Twenty milliseconds before a laser shot the unit is triggered and scans all the enabled inputs to obtain a background reading. About 45 ms after the shot, when EM interference due to the shot has subsided, the enabled inputs are again scanned for a period of 4 seconds to allow the peak reading of each input to be obtained. The background reading is subtracted from the peak and the unit may then be read via the CAMAC dataway by the GEC 4080 computer where the readings are converted to joules, using appropriate calibration factors, for display and logging.

1.3.4(d) Wave Plate Attenuators

During the recent laser upgrade variable optical attenuators have been installed and the operating procedure changed. Previously laser output was controlled by adjusting amplifier flash lamp voltage settings but now the voltages are fixed while laser output is controlled using the attenuators. Attenuation is varied using rotating waveplates (section 1.3.1(b)) driven by stepper motors which rotate the waveplate by 0.1° per step. A unit has been constructed in a EUROCARD support frame incorporating an "Essex Tiny Basic Computer" and driver cards for up to 4 stepper motors. A program written in BASIC resides in PROM in this micro-computer which issues step pulses to the motor drives, keeps track of motor position, receives commands from the front panel push buttons and communicates with the GEC 4080 via an RS 232 unit in the CAMAC dataway.

M Forster, C Reason and P Gottfeldt (RAL)

1.3.5 Extended Clean Room Facility

In order to accommodate the assembly of three new disc amplifiers and regular servicing of nine disc amplifiers an extended clean room area has been established. This extended area includes several new facilities, essential for testing, cleaning and assembly of components associated with the laser system.

The principal facility added to the clean room is the high pressure freon spray booth. This consists of a sealed box capable of holding the component parts of a disc amplifier which are then cleaned using a hand held gun spraying a freon-alcohol mixture of approximately 1000 psi. The freon-alcohol mixture, after spraying, is collected, distilled and stored in a large tank ready for re-use. The cleaning of components, using this method, has proved both faster and more hygienic than the previous portable 80 psi freon spray.

Another facility added to the clean room area is the disc amplifier flash lamp test facility. This facility allows the carrying out of infant mortality tests on new flashlamps now being manufactured in this country, prior to installation in disc amplifiers (see 1.3.7).

A CW YAG laser has also been installed in the new clean room area. This laser is used to carry out acceptance tests on various optical elements, including the determination of the transmission and reflection properties of dielectric mirrors and polarisers.

The Zygo interferometer and its associated vibration isolated table have also been installed in the clean area, enabling most tests on optical elements to be carried out under clean room conditions.

Additional storage facilities have also been installed, in order to accommodate the wide range of spares required for the disc amplifier and

the many optical spares for the remainder of the system.

N Allen, S Sunderland and R Wellstood (RAL)

1.3.6 Single Beam Disc Amplifier

Future upgrade proposals for the glass laser system include an increase in performance of the single beam line to meet future needs for single beam experiments and to keep the backlighting beam power in line with advances in power and energy available for compression experiments. The first stage of this single beam upgrade would provide twice the energy and power currently available with a ϕ 108 mm disc amplifier and would therefore need to be of about twice the beam area. Thus it has been decided to build a prototype ϕ 150 mm amplifier in a form as efficient and economic as possible. Figure 1.11 shows the design of this amplifier. The 'box' design gives high efficiency and fewer lamp circuits; larger bore lamps (ϕ 20 mm 1/D as opposed to ϕ 15 mm 1/D) improve efficiency and reduce the number of lamp circuits; and optimum concentration of neodymium at 3% gives high gain and hence low B-integral. Using LLNL gain data for the ϕ 150 mm amplifier should give a small signal gain of nearly 5 using only 12 lamps (6 lamp circuits) and a stored energy (180 kJ) less than currently used to drive the ϕ 108 mm amplifiers.

Figure 1.12 gives the expected performance of a single beam line using this amplifier showing several modes of operation. A prototype amplifier and power supply is currently under construction.

I N Ross (RAL)

1.3.7 Flashlamp Test Facility

In operation flashlamps used in the disc amplifiers are conservatively run and in principle a life of $\sim 10^6$ shots can be anticipated. However, due to presently unidentifiable defects, some lamps exhibit a so called "infant mortality" and do not survive the first few shots. Until recently the lamps used have been subjected to an infant mortality test at source

but a new UK supplier has been unable to provide this service. The availability of a test facility at RAL was imperative in order to eliminate lamps showing infant mortality before installation in the disc amplifiers.

The test procedures are based on LLNL specification LES 2292. The lamps which are 20 mm diameter and 1295 mm between electrodes are tested singly at twice the normal energy, ie 20 kJ. The pulse forming network was designed to produce a 10% to 90% rise time of 150 μ s, a peak current of 7 kA and to be critically damped. Details of the PFN together with a typical current waveform are shown in Figure 1.13. It can be seen that the design parameters have been substantially met apart from slight underdamping.

The facility is designed to test lamps semi automatically. It consists of four units; a charging supply; the PFN, ignitron and lamp housing (see Figure 1.14); a local controller and an Apple II microcomputer. The charging supply is manufactured by Hartley Measurements Limited and is capable of charging at 600 J/s. The capacitor and lamp housing is in two sections. The lower section contains the PFN, ignitron and dump circuit while the upper section houses the flash lamp mounting and pulse current transformer. The Apple II microcomputer is isolated from the high voltage and electrically noisy environment of the pulse circuits by a fibre optic link to the local controller based on a 6502 micro-processor.

The local controller directly controls the dump relay, the power supply and the ignitron firing but only under command from the Apple, the only action it can initiate on its own is dropping the dump relay and turning off the power supply. It will do this if the dump relay is up and it has not received a command from the Apple in the previous 50 ms.

As explained the Apple does the controlling, data gathering, analysis and storage. It does this with a 10 K basic program supplemented by 1 K of machine code and $\frac{1}{2}$ K of shape table. The program is menu based for ease of use and leads the operator through the necessary procedures until it starts the charge sequence. From then on the facility is completely

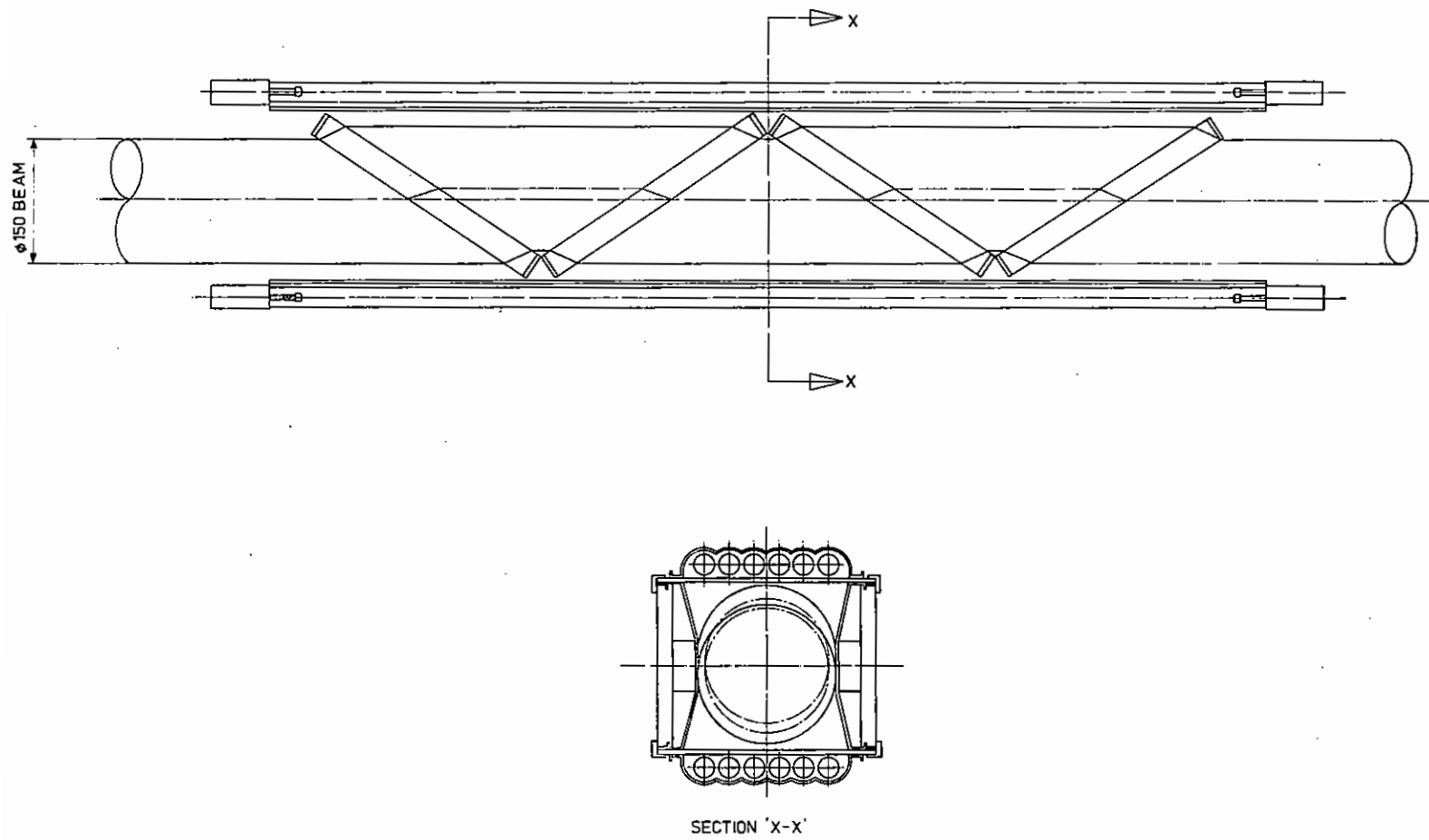
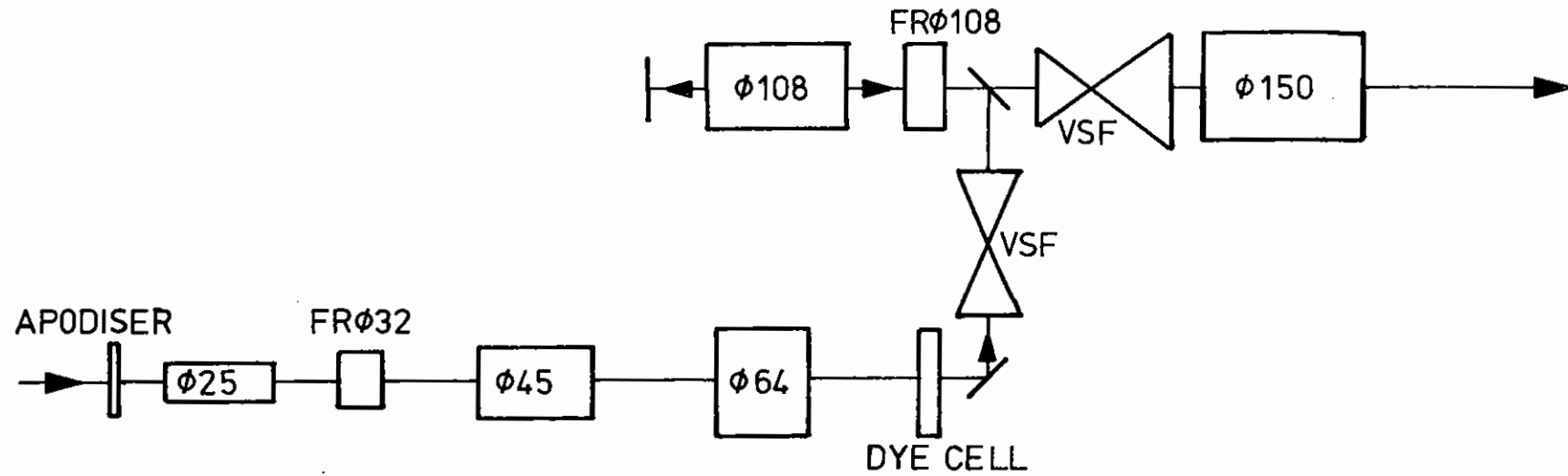


Fig. 1.11 Design details of the Ø150 disc amplifier.



PULSE LENGTH	ROD CHAIN	$\phi 108$ SECTION	$\phi 150$ SECTION	FILL FACTOR
0.1ns	E = 3.6J $\Delta B = 1.53$	E = 58J $\Delta B = 3.9$	E = 163J $\Delta B = 2.75$	0.65
0.3ns	E = 78J $\Delta B = 1.11$	E = 114J $\Delta B = 2.70$	E = 300J $\Delta B = 1.70$	0.67
1.0ns	E = 19J $\Delta B = 1.15$	E = 199J $\Delta B = 1.7$	E = 452J $\Delta B = 0.82$	0.70

Fig. 1.12 Expected performance of beam line using Ø150 disc amplifier.

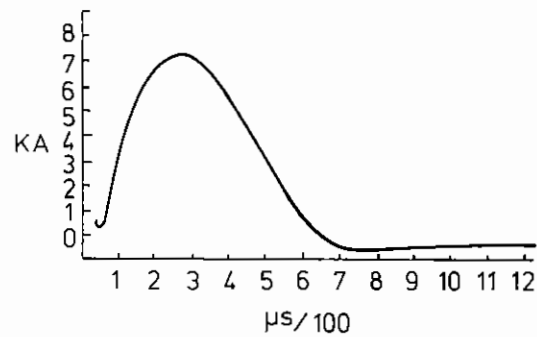
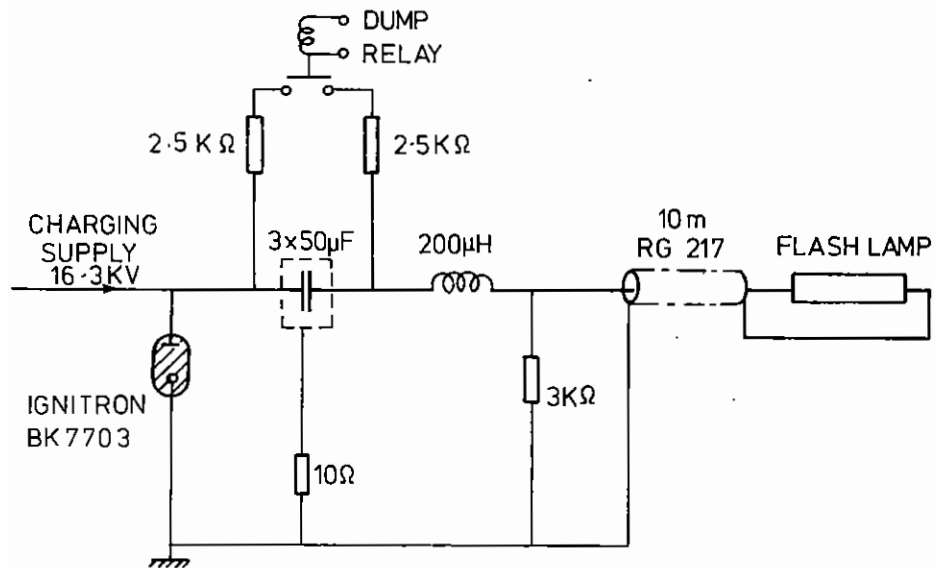


Fig. 1.13 a) Details of the pulse forming network.
b) Typical current waveform.

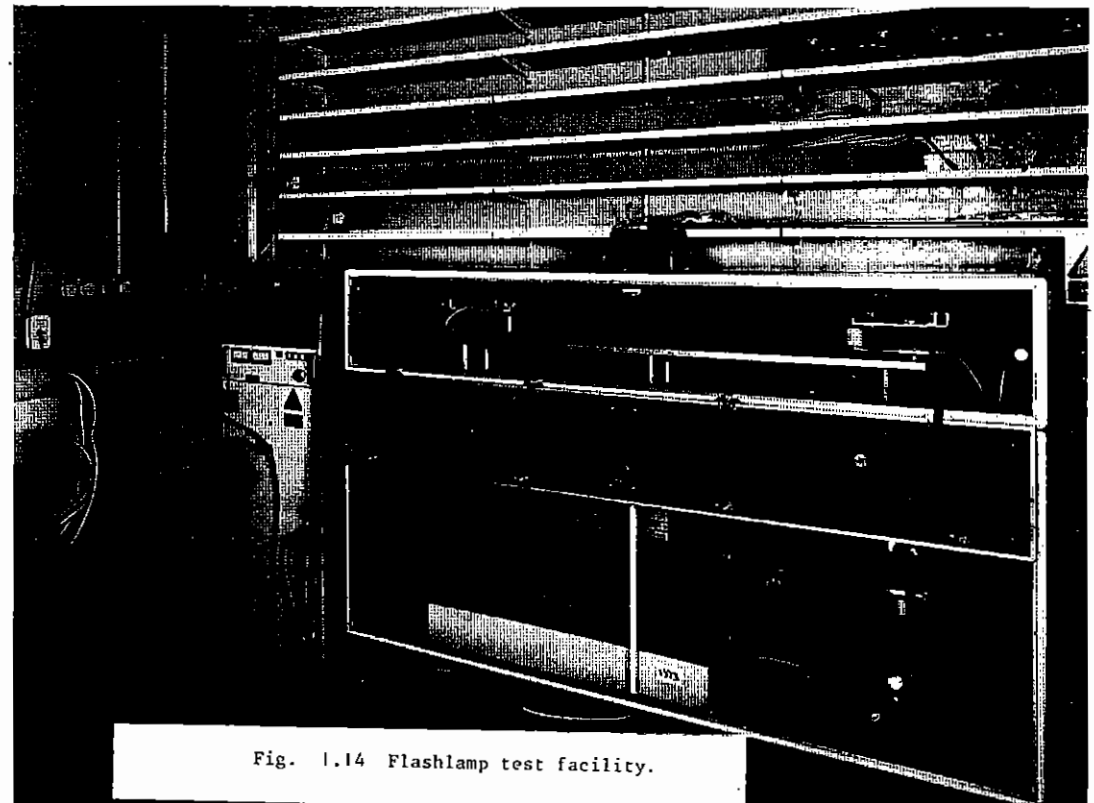


Fig. 1.14 Flashlamp test facility.

automatic until the test is complete or a lamp fails. The program also contains sections to recover and display data on disc and to test the electronics.

After each shot the Apple collects the current waveform from a digitizer in the local controller, displays it on the screen and then stores it on disc. A patch into the Apple's disc operating system was written to do this efficiently. The peak current and time to peak current are then measured and if both are within preset boundaries the lamp is deemed to be good and if there are still shots to be done the sequence is restarted.

Lamps have been observed to fail in a number of ways, some continuing the discharge almost to completion before showing any perturbation of the normal current waveform. However, it does not seem that these waveforms can shed any light on the precise mechanism of lamp failure.

M Forster, R Day and P Gottfeldt (RAL)

1.4 Target Areas

1.4.1 Existing Target Areas

In order to correct an imbalance in demand for the two target areas caused by user requests for X-ray backlighting experiments in plane geometry in addition to the spherically symmetrical experiments for which Target Area 1 was designed, a second beam was brought into Target Area 2. It is derived from the 6 beam driver disc amplifier following the output vacuum spatial filter. As with the original Target Area 2 beam derived from the "backlighting" channel, one remote controlled mirror can be inserted to switch the beam from one target area to the other. Either or both beams can be delivered to Target Area 2 with independent long or short pulse operation. Either beam can be brought to the north, south or top parts of the target chamber: the layout and wavelength combinations available are shown in Figure 1.15. An additional remote controlled target mount has also been installed for backlighting experiments.

The alignment systems in Target Area 2 have been improved to permit more rapid wavelength changes. All-mirror optics are now used in TV comparators, and Cassegrain telescopes are now used to observe the target as well as to produce the expanded alignment beams. The injection of CW lasers at 1.064 μm , 0.633 μm , 0.514 (or 0.528) μm and 0.351 μm has been simplified for the target area operator. The layout is shown in Figure 1.16.

An improved method of harmonic crystal alignment is now in operation in Target Area 1. This uses a wild self-levelling level as an absolute horizontal reference to ensure that the laser beams incident on the harmonic conversion crystals can always be maintained at the same angle to the crystal e-axes, which are vertical.

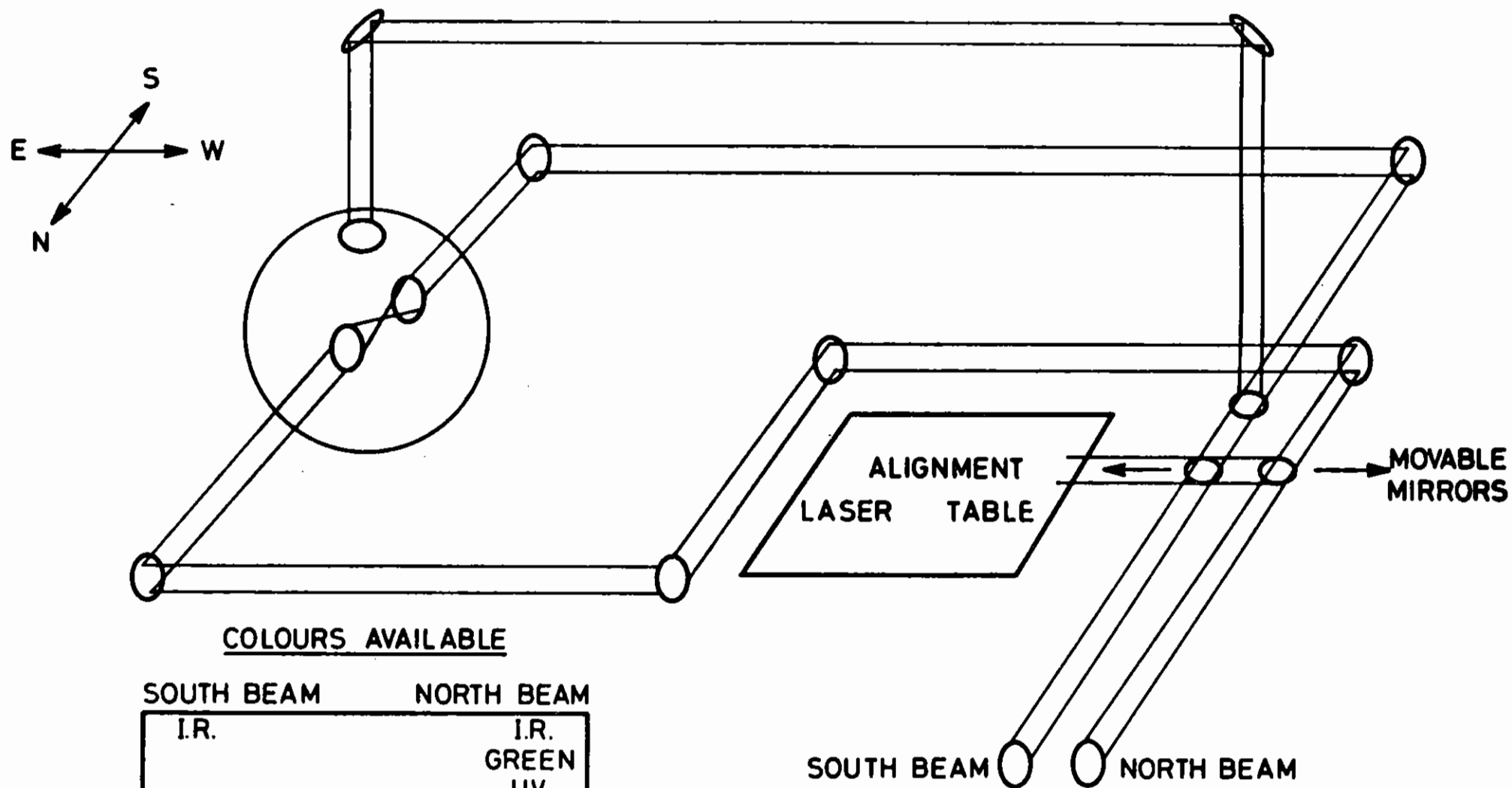
R Eason and W T Toner (RAL)

1.4.2 New Target Areas

During the past year detailed design studies have been carried out to establish the space and equipment required for the construction of a 12 beam target area. These have resulted in the construction of a building extension measuring 25 metres x 12 metres on the north side of the existing target areas.

Contracts have been placed for the completion of this building, and modifications to Target Area 2, so that by the Autumn of 1983, two new target areas will be available for fitting out (Target Area East and Target Area West) together with a central control room and air conditioning plant room. Access to these areas is from the existing control room via a mezzanine floor in Target Area 2, which will also eventually provide space for a beam switch yard. Target Area West is in line with Target Area 1 and will house the 12 beam target facility, the switch yard being used to divert the beams into Target Area West.

Contracts have also been placed for the construction of a new spherical



COLOURS AVAILABLE

SOUTH BEAM	NORTH BEAM
I.R.	I.R. GREEN U.V.
GREEN	I.R. GREEN U.V.
U.V.	I.R. GREEN

Fig. 1.15 TAI Upgrade layout.

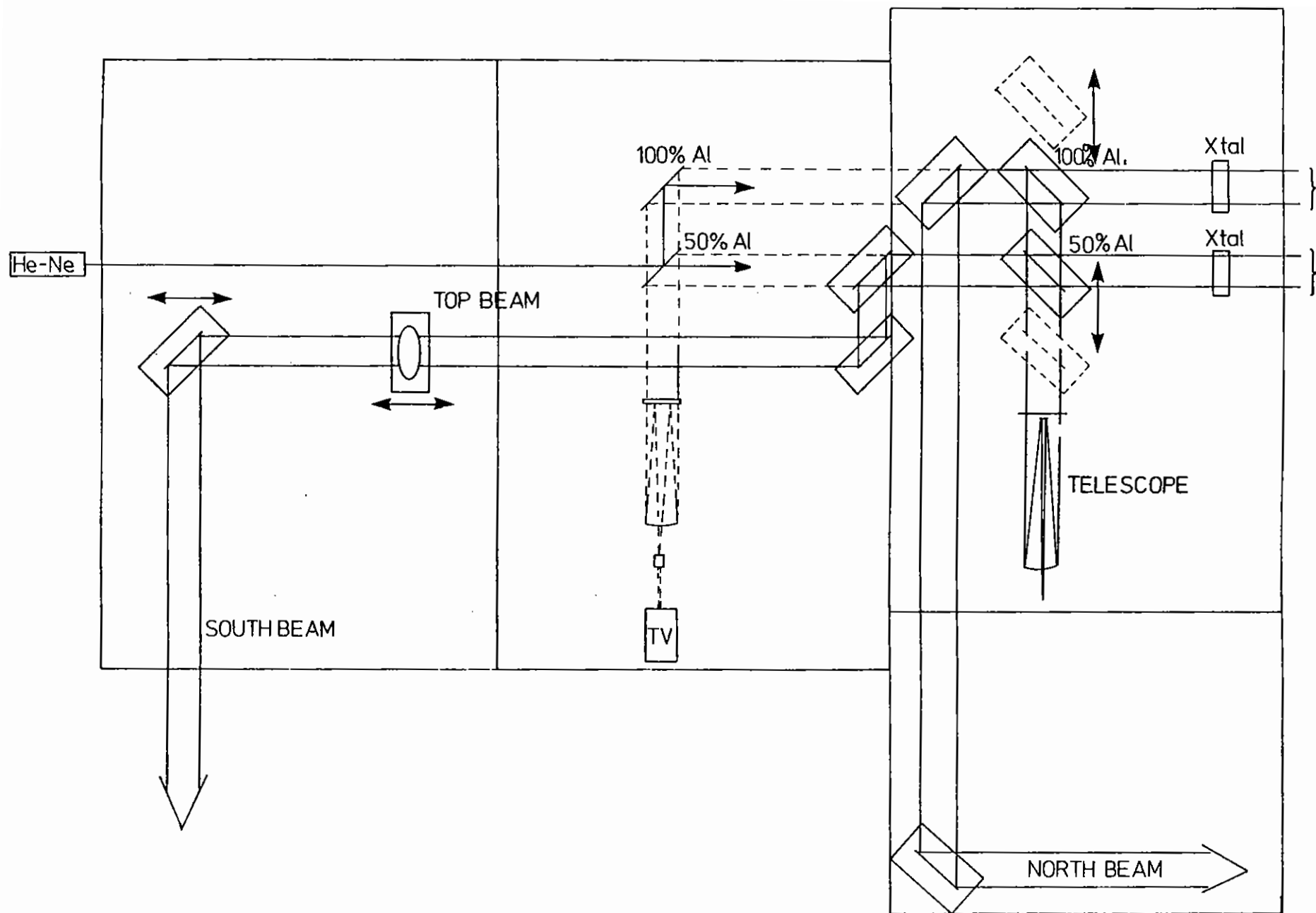


Fig. 1.16 TAIL 2 beam alignment system.

vacuum chamber for the twelve beam system. This is scheduled for completion by September 1983.

J E Boon and W T Toner (RAL)

1.5 Optical Development

1.5.1 Aspheric Lenses

A considerable amount of work has been carried out on the design and manufacturing procedure of lenses for the proposed twelve beam system. The final F number of the focusing optics was chosen on the basis of several conflicting requirements but paramount amongst these was the requirement for uniformity of illumination based on our theoretical estimates and outlined in last year's Annual Report.

Two lens systems have been designed based on a focal aperture of F/2.5. The first design is a doublet lens in fused silica based on previous design (1.06). It is found in practice, with proper balance of power between components, that a lens covering the wavelength range $250 < \lambda < 1054$ nm with virtually diffraction limited imaging can be readily designed. Such lenses have a reflective loss of ~ 20% in the absence of antireflective surface coatings. In practice it is found that such coatings damage at high power levels, in most cases, forcing the lenses to be operated uncoated. A singlet lens, of suitable optical quality and able to operate over the wavelength range of interest, would offer a considerable advance over lenses in current use.

The design of such a multiwavelength singlet lens has been carried out (1.07). The principle underlying the design of the lens is that, unlike the conventional aspheric, which is made to satisfy Fermat's principle at one particular wavelength, the lens is designed with an aspheric surface that is perturbed so as to minimise the aberration at more than one selected wavelength. This can be done by balancing the stigmatic performance of the lens at its nominal, or design wavelength, against the spherochromatic aberration incurred in operation at some other wavelengths

by successively adjusting the coefficients defining the aspheric surface. This procedure should not be confused with the more common damped least squares minimising procedure operating on the Seidel coefficients (which is in fact a particular case of this more generalised case).

Theoretical investigations suggest that a lens based on this principle and covering the entire wavelength range of interest could be designed whose performance would be such that it would produce the necessary uniformity of illumination when used in a twelve beam system. A trial lens is currently in the process of being manufactured.

D J Nicholas (RAL)

1.5.2 Uniformity of Multi-Beam Illumination

1.5.2(a) Uniformity for Six Beam Illumination

Calculations have been carried out with respect to the parameters which affect implosion uniformity (1.06)(1.07). In particular it was useful to look for the best conditions for uniformity with six beam illumination, where the adjustable parameters are the input beam profile (assumed to be a supergaussian with variable power index), the input beam diameter, the lens focal length and the target defocus distance. For an assumed inverse square electron density distribution at the target surface the energy deposition in the corona is given by:

$$D(\phi) = I(\phi) \cos^2 \phi$$

where ϕ = ray angle of incidence at target surface and I = incident intensity.

Using this relation the non-uniformity of energy deposition has been calculated for six beam illumination. Figure 1.17 shows parameter regions giving less than 5% nonuniformity where the target defocus = p x target radius. It can be seen that if N is chosen to be say 1.3 (not much 'squarer' than a gaussian) and the beam width parameter to be say 20° then

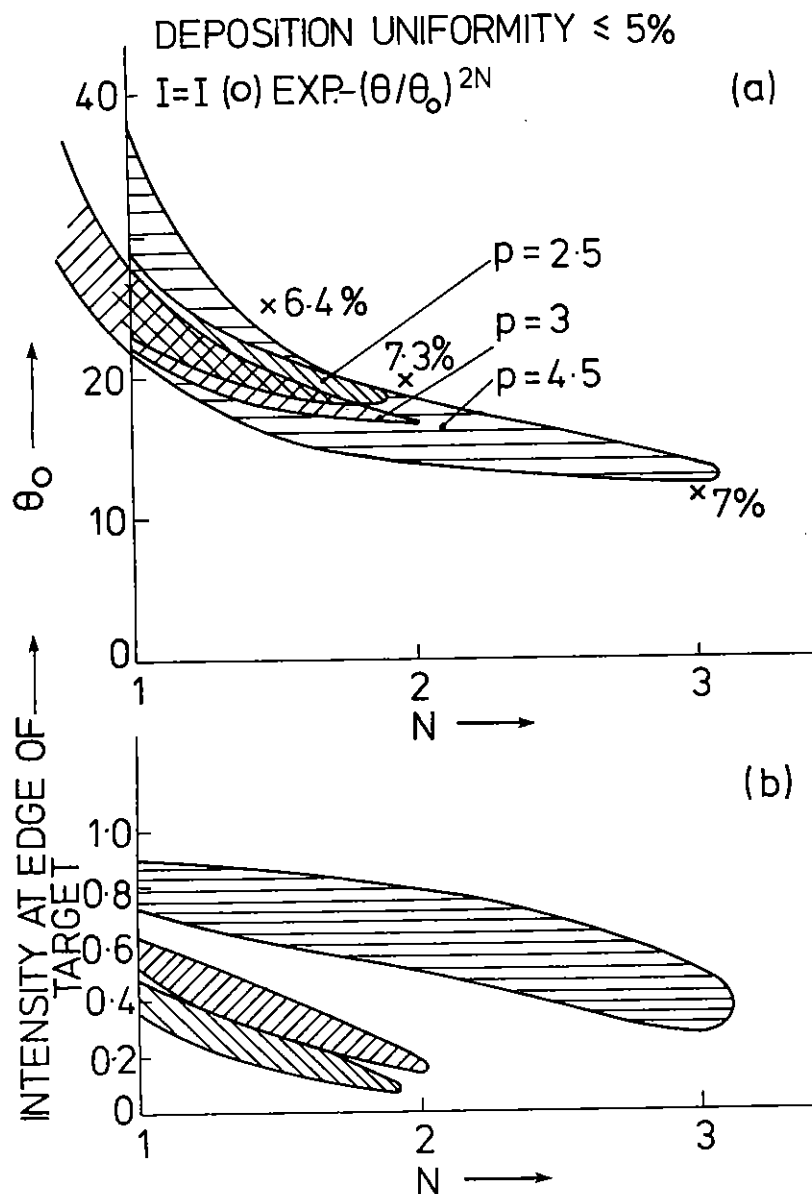


Fig. 1.17 Deposition uniformity with 6-beam illumination a) parameter regions giving less than 5% non uniformity: spot values refer to $p=4.5$ (b) corresponding values of the beam intensity at the edge of the target.

good uniformity can be achieved with a reasonable tolerance to changes in the parameters. Also shown are the corresponding values of intensity for each beam at the edge of the target, giving an indication of the extent to which energy missed the target.

Figure 1.17 applies early in the laser pulse when the plasma scale length is still small compared to the target size. As the plasma expands so the deposition moves outwards resulting in improved thermal smoothing. Also the energy in each incident ray becomes deposited over a considerable distance along the ray path (over a considerable length across the target surface) leading to a second smoothing process, refractive smoothing. An illustration of the deposition non-uniformity later in the laser pulse when the plasma scale length is approximately equal to the target radius is shown in Figure 1.18. The effect of refractive smoothing in giving good uniformity over a wider range of parameters is evident by comparison with Figure 1.17. It can also be seen that good uniformity can still be achieved later in the pulse with the same beam parameters.

I N Ross (RAL)

1.5.2(b) Uniformity of Multi-Beam Illumination

In last year's CLF annual report we described calculations of the uniformity of illumination of a spherical target with 6 and 12 super-Gaussian laser beams. The calculations did not address the problem of absorption. The uniformity of a spherical implosion is determined by the uniformity of drive pressure rather than the uniformity of illumination. The generation of pressure is affected by the fractional absorption of the incident laser light and by the efficiency of its transport from the region of absorption in the corona to the dense plasma at the ablation surface. The relationship between the incident energy and the drive pressure is complicated. It depends on such factors as the density profile in the corona, which will vary during the course of an implosion, and the ratio of resonance to inverse Bremsstrahlung absorption. Consequently, we do not attempt to represent this relationship accurately, but instead explore the sensitivity of the drive pressure uniformity to

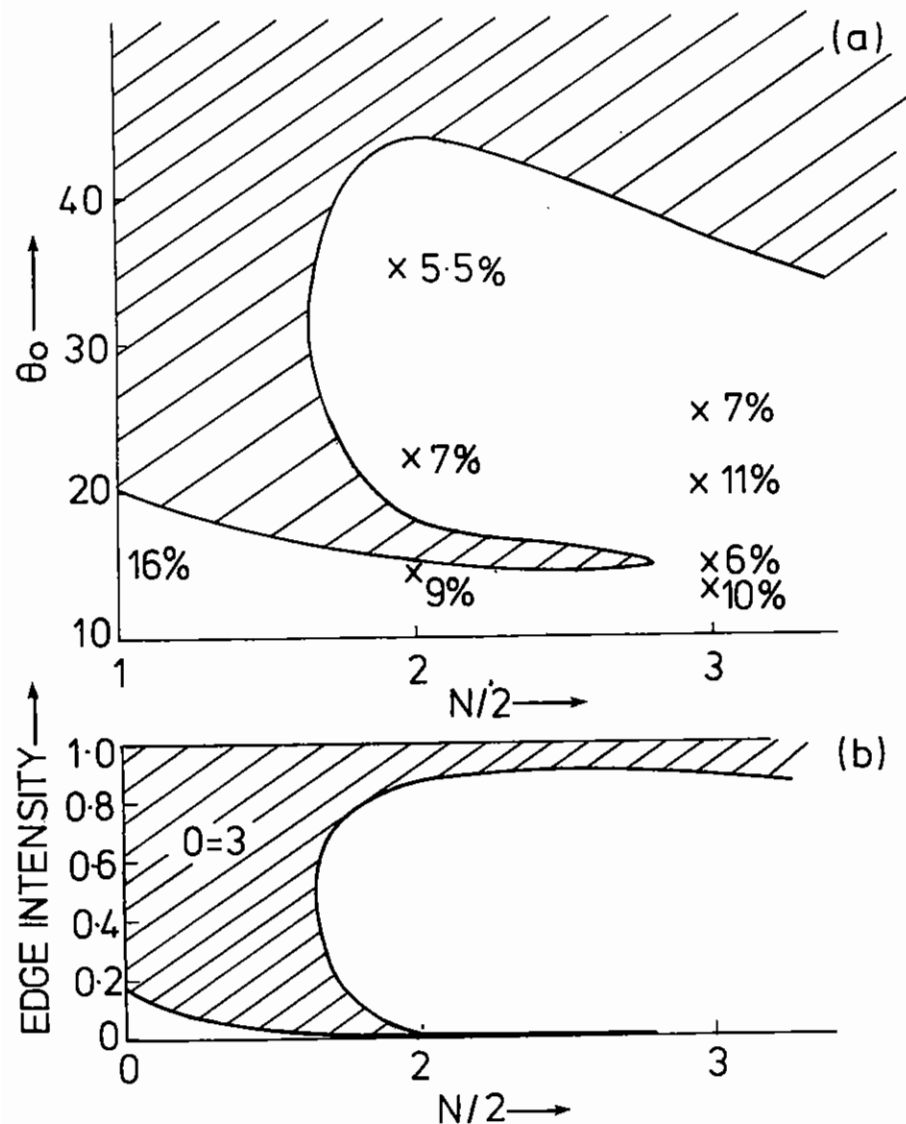


Fig. 1.18 Uniformity of ablation pressure with six beam illumination and long scale length plasma.
 a) Parameter regions giving less than 5% non uniformity
 b) Corresponding values of beam intensity at the edge of the beam.

this relationship. We assume that the drive pressure P generated by each beam at one point on the target is related to the incident irradiation I by

$$P^\alpha = I (\cos i)^n$$

where i is the angle to the normal at which the laser radiation is incident on the target. We then calculate the pressure uniformity for values of $0 \leq n \leq 5$. If the drive pressure is proportional to the incident irradiation then the correct value for n is zero. If the energy is weakly absorbed on a linear density ramp then $n = 5$ (1.07). We can expect n to lie somewhere within this range.

The percentage non-uniformity, defined as $100 (P_{\max} - P_{\min}) / P_{\min}$, is shown in Figures 1.19 and 1.20 for second, third and fourth power super-Gaussian beams focused at a distance s times the target radius beyond the target centre. The value of n is different for each graph. $n = 0, 1 \dots 5$ for graphs a, b ... f respectively. Figure 1.19 is a plot of the results for 6 f/l lenses, and Figure 1.20 is a plot for 12 f/2.5 lenses. The 12 beam system is found to be much less sensitive to the value of n , as might be expected, because the beams are more densely packed on the surface of the target so that rays at large angles of incidence are less important. The poor uniformity found in some cases such as $n = 4, 5$ for 6 lenses could in principle be improved by boosting the outer part of the beam but this requires a knowledge of the actual relationship between P and I .

A R Bell and D J Nicholas (RAL)

1.6 Target Production and Development

The routine preparation, characterisation and mounting of microballoon, microdisc and multilayer foil targets has once again occupied most of the time of the target fabrication group during the last year.

Target development activities have centred around investigations of new methods for producing structured microdisc and microballoon targets for

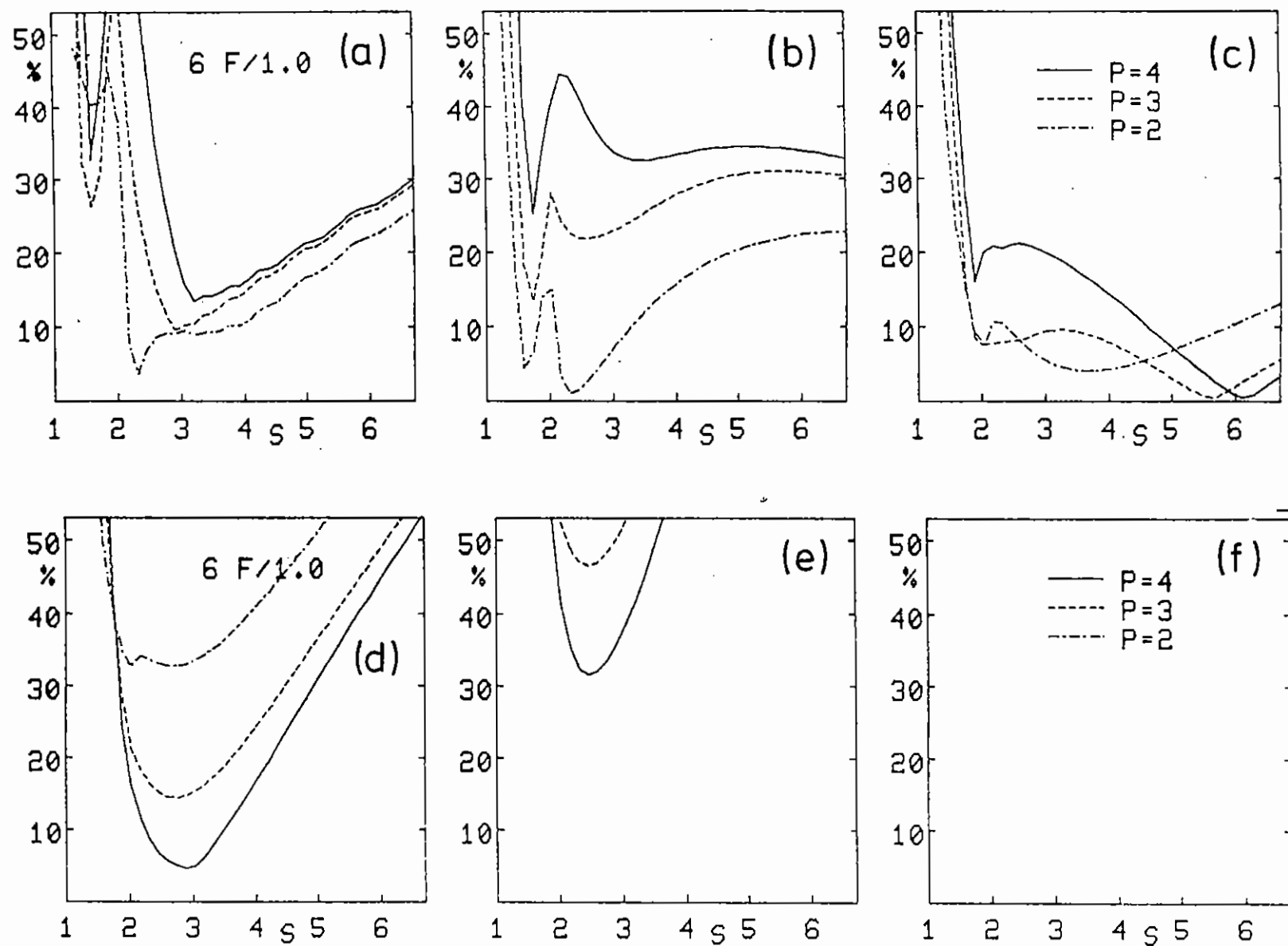


Fig. 1.19 Percentage non-uniformity for 6 super-gaussian beams with super-gaussian power $P = 2, 3, 4$ focussed by $F:1$ lenses. S is the position of the focus in units of the target radius, beyond the centre of the target. The cosine absorption index is $n = 0, 1, 2, \dots, 5$ in Fig. 1.19a) 1.19b) 1.19f) respectively.

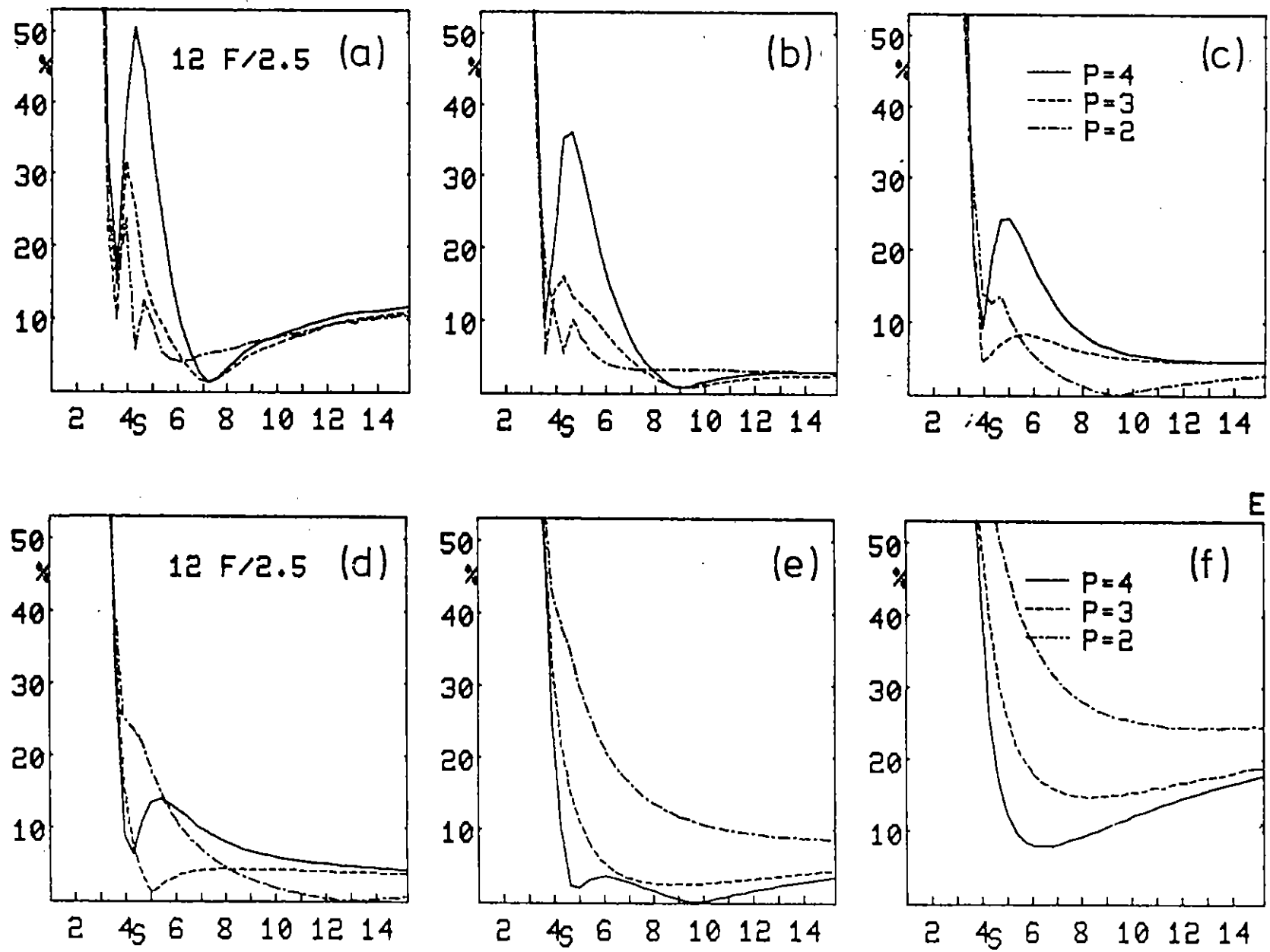


Fig. 1.20 Same as Fig. 1.19 but for 12 F/2.5 lenses.

experiments studying growth of Rayleigh-Taylor instabilities. Preliminary experiments of this type using streaked X-radiographic techniques to observe the instability growth rate in accelerating targets have already been performed successfully using rippled microdisc targets. These targets have been produced relatively easily by evaporation of 2-3 μm of aluminium through a mask, having an array of 500 μm diameter apertures, onto the surface of a Moiré fringe grating of suitable line periodicity. Gratings available have deposited chrome stripes 0.5 μm high of periodicity from 150 to 20 lines per mm (ie stripe/space widths from 5 μm to 25 μm). The evaporated microdiscs produced by this method replicate the grating surface structure well and when released make purely rippled targets with minimal variations in areal density across the target surface. Such targets are ideal for transmission X-radiography experiments.

Two new target structuring processes have been investigated. The first of these involves reactive ion beam milling and consists of irradiating the surface of a polymer with a beam of 5 keV oxygen ions from a small ion source. At current densities of about 10 mA/cm² etching of the surface at a rate of a few microns per hour occurs and consequently if a metallic mask is placed in front of the polymer the pattern of the mask is etched onto the surface by the collimated ion beam. Etching rates of metals by oxygen ions has been found to be significantly lower than polymers and consequently mask erosion is not a problem. Various nickel masks have been made, by photolithography techniques, consisting of bars and spaces in a comb like structure with bar and space width of 10, 20 and 40 μm . Using these masks line patterns have been milled successfully up to several microns deep in the surface of mylar microdisc targets. Such targets are not as suitable as rippled microdiscs for transmission X-radiography experiments due to the variation in areal density across the target surface. They are however suitable for X-ray backlighting observations of Rayleigh-Taylor instability growth though exact alignment of the microdisc ripples and observation axis is difficult. To overcome this problem the ion milling technique has been extended to structure one side of mounted polymer coated microballoon targets of 200 -300 μm diameter. Due to their curvature such targets are ideal for observation

of instability growth at the ablation surface when irradiated on the structured face using 3 beams of the 6 beam compression facility and using pulsed X-ray backlighting methods. Figure 1.21(a) shows a SEM photograph of a polyxylyene coated glass shell with a 10 μm bar pattern produced by ion milling.

The second target structuring method is novel and relatively simple and utilizes a small rare gas halide laser. It has been found that when polymers with high absorption coefficient are irradiated with ultra-violet photons from either ArF ($\lambda = 193 \text{ nm}$) or KrF ($\lambda = 247 \text{ nm}$) lasers at pulsed power densities of about 10 MW/cm² decomposition of a thin surface layer occurs and material is removed. At an energy density of about 0.1 J/cm² about 30 nm of material is removed per laser pulse (15 ns) from mylar (PET). By placing a metallic mask in close contact with the polymer suitable patterns can be photo-etched into the surface. Since this process is so simple and can be carried out in air with no chemical development phase it is of great interest at present to determine its relevance to integrated circuit lithography processes (see section 3.6) .

Using this laser etching technique with the 10 μm bar and space nickel masks, acquired for ion beam milling, patterns up to a few microns deep have been etched very successfully in the surface of 12 and 25 μm thick mylar foils. 550 and 270 μm diameter microdisc targets have been prepared from these foils using precision micro-punches. Figure 1.21(b) shows such a laser etched microdisc target. Work is continuing to further develop and understand the process and to determine how etch rate depends on photon energy, power density, polymer structure and absorption coefficient.

The manufacture of pure polystyrene shells by solvent drying techniques has been described in detail in last years report. The problem of such shells for X-ray spectroscopy experiments is that gas retention lifetimes are very short for suitable gases such as neon and argon. Attempts to manufacture shells by similar techniques from polymers with lower permeabilities have so far been unsuccessful and consequently alternative methods of solving the poor gas retention problem of polystyrene shells have been evolved. Various high molecular weight gases containing

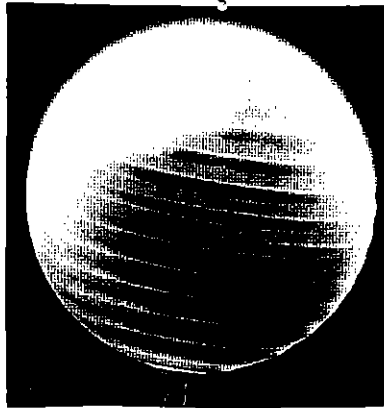


Fig. 1.21 (a) S.E.M. Photograph of a polyxylene coated glass shell target with a $10\ \mu\text{m}$ pattern produced by reactive ion milling.

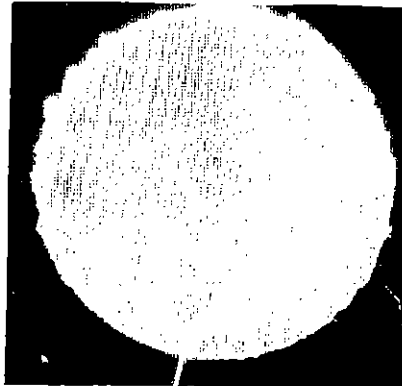


Fig. 1.21 (b) S.E.M. Photograph of an excimer laser etched microdisc target.

suitable high atomic number constituents have been investigated and their polystyrene permeabilities checked. Measurements have indicated that Freon 12 (dichlorodifluoroethane) has a sufficiently low permeability for experimental purposes and that when polystyrene shells of $200\ \mu\text{m}$ diameter and $10\ \mu\text{m}$ wall thickness are immersed in Freon 12 at room temperature the air diffuses out and is replaced by the Freon in roughly one hour. This time is not long enough to store Freon filled target for lengthy periods but is adequate for implosion experiments if targets are transferred rapidly from the filling rig to the target chamber. Implosions in which core densities have been diagnosed by X-ray spectroscopy of chlorine line emission have been performed successfully using Freon 12 filled polystyrene shells. In these experiments delay times between removal of the target from the filling rig to the time of the laser shot were reduced to less than 5 minutes during which period fill gas pressures would fall by only a few per cent.

C Brown, B Child, S Halletwell, K Hollingworth, N Prior and P Rumsby (RAL)

1.7.1 X-ray Framing Camera

X-ray and optical streak cameras such as the Hadland 675 can be used in framing mode, but repeated frames can only be obtained at about one frame every two nanoseconds. To be of interest in laser compression experiments a framing rate of ten times this value is desirable although useful information could be obtained with slightly slower framing rates.

In a conventional streak camera operating in framing mode, a staircase voltage is applied to the deflection plates so that an image is moved successively to different points across the screen. The electron beam may be blanked off during the movement but in many cases this is not necessary since exposure in between the image positions is relatively low. The major limitation to the framing rate of such a streak camera arises because of transit time of electrons past the deflection plates. Thus even if one were able to apply a very short rise time pulse to the deflection plates, the image would not completely have moved until one electron transit time later. This time is typically 400 ps. In addition

to this, capacitance deflection plates are difficult to drive with very fast rise time pulses.

Both of these problems are overcome by using a travelling wave deflection structure whose effective transit time can be reduced to 50 ps and which can be readily impedance matched to strip line, microstrip or coaxial cable driving circuitry.

Various configurations for the travelling wave structure have been considered and one similar to that shown in Figure 1.22 is based on the EEV P855 tube but has been built in a metal demountable form to allow ease of modification.

Electrical rise time measurements on the deflection structure have shown a 10% - 90% rise time of about 70 ps. If this rise time were translated directly to the electron beam it would allow framing times of approximately 200 ps whilst still maintaining an acceptable resolution.

Static image resolution tests have shown a cathode resolution of about 18 lpp mm. This can be maintained with only a small loss of resolution over the 50 mm diameter screen. Deflection sensitivity is about 30 volts/mm.

The fast rise time staircase waveform required to drive the camera is produced by a silicon photo-conducting switch (Auston switch). Although the silicon switch has to be used with a pulsed bias (to avoid switch breakdown) it has a long carrier lifetime (> 100 ns) which is ideal for this application. Thus when the switch is illuminated by a series of very short laser pulses, which individually are incapable of driving the switch into saturation, the silicon conductivity increases in a series of steps producing a fast rise time staircase waveform. In the early experiments the laser pulse train has been produced from a single short laser pulse by means of a Fabry Perot but this results in an output waveform with rapidly decreasing step height and tends to result in later frames overlapping. Various schemes are under consideration to overcome this problem.

In the initial dynamic trials a 1.5 ns 10 J pulse was focused onto a brass

target to act as a soft X-ray source for the framing camera. Two grids of spacings 0.1 mm and 0.4 mm were placed directly in front of the gold coated beryllium cathode to act as simple resolution test. A synchronised 5 mJ 100 ps probe pulse was used to illuminate the silicon switch after passing through the Fabry Perot. The interpulse time was determined by the Fabry Perot spacing.

Framed images of the grids were obtained for two settings of the Fabry Perot producing 420 psec and 300 psec interframe times. The image resolution was degraded along the time shift direction. At 420 psec interframe time, the time axis photocathode resolution was 6 lppmm whereas at 300 ps frame time the resolution had dropped to about 3 lppmm. This resolution would be consistent with a step rise time of 120-150 ps. since the laser probe pulse width was ~ 100 ps this suggests that the camera deflection system has a rise time of 65 - 110 ps. This is not inconsistent with TDR measurements of the electrical rise time.

T A Hall, C Kler and E McGoldrick (Essex University)

1.7.2 Sub-nanosecond Gated X-ray Intensifier

A gated X-ray intensifier with a resolution of 50 ps has been developed by a collaboration of Imperial College, AWRE Aldermaston and RAL.

Figure 1.23 shows the device. The body, housing a copper mesh, microchannel plate, phosphor and fibre optic output window is constructed from ceramic tubes and rings. The sections are vacuum-brazed together. The microchannel plate is a Varian, type VW8900Y, with a pitch of 14 μ m and operating voltage up to 1 kV with a gain of $2 \cdot 10^4$. The phosphor is a P11-type operating at 5 kV and having an opaque aluminium layer over it. The fibre optic window has 5 μ m diameter fibres. The copper mesh has a pitch of 25 μ m and transparency of 70%. This provides a ground plane for the 50 Ω strip transmission line in the vicinity of the photocathode. The 500 nm photocathode is made of Caesium Iodide, vacuum deposited onto aluminised Mylar (3 μ m Mylar, 100 nm Aluminium). The Aluminium layer is required to carry the gate pulse current and provide a conducting plane to

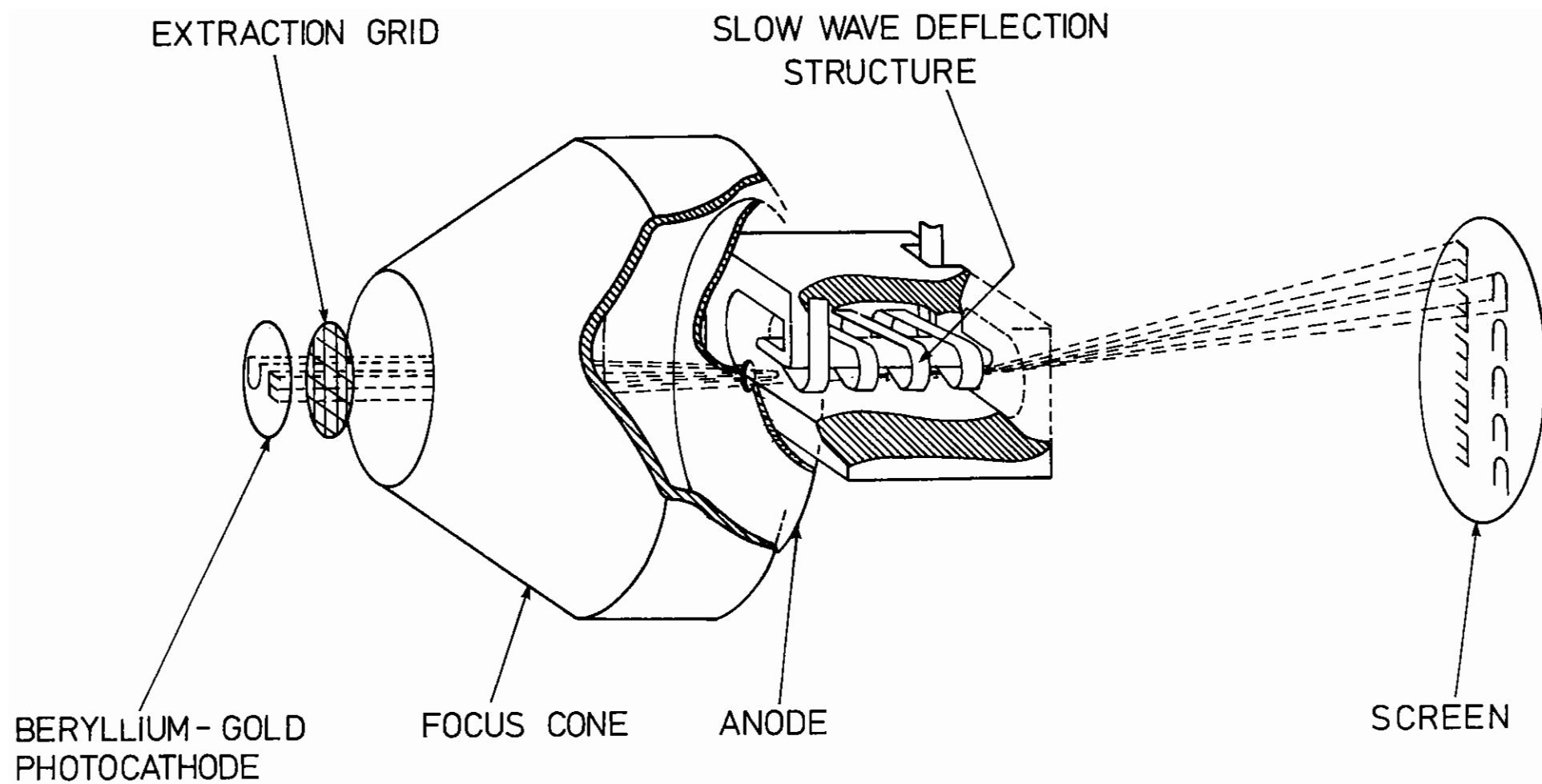


Fig. 1.22 Schematic of X-ray framing camera.

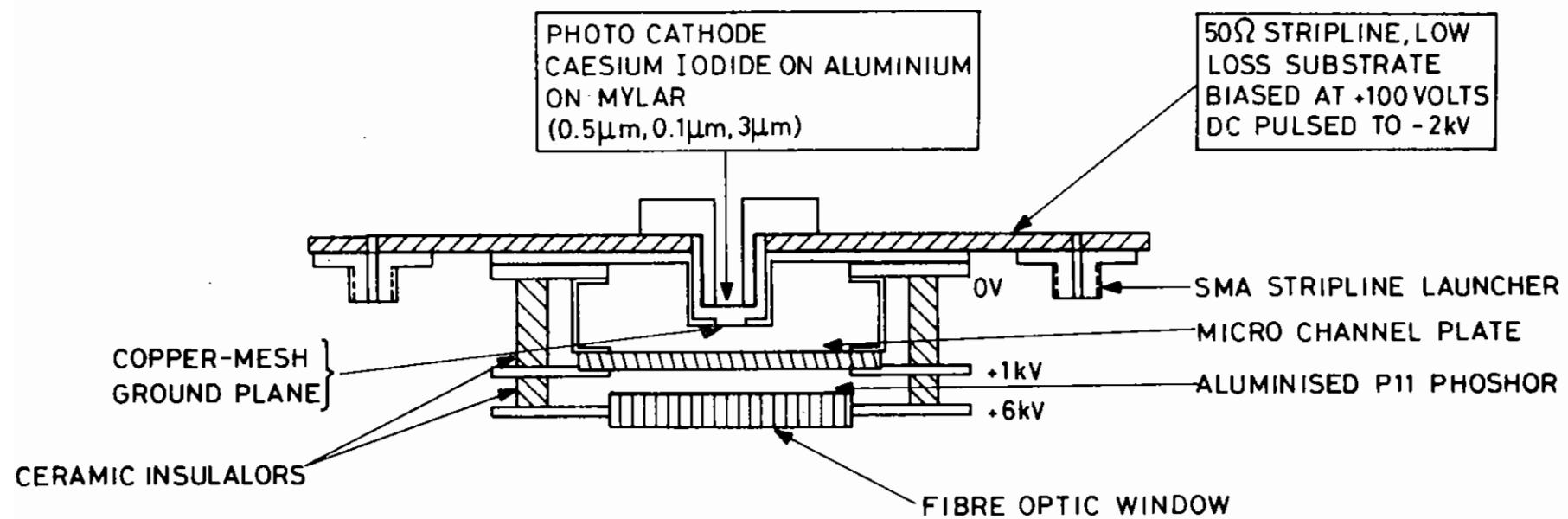


Fig. 1.23 Schematic of the gated intensifier.

ensure that the electric field is perpendicular to the plane of the photocathode. The output from the device is recorded on HP5 film discs.

The gate pulse was produced with an Auston switch using a Gallium Arsenide (GaAs) switching element which can maintain a DC electric field greater than 1 kV mm^{-1} . The GaAs was 6.5 mm wide, 10 mm and 0.3 mm thick. Electrodes 4 mm apart were painted on with a silver preparation which was also used to stick the chip across a gap in a strip transmission line. The two-strip transmission lines, the Auston switch and the photocathode assembly, were connected with 6.35 mm diameter semi-rigid co-axial cable. This provided for remote positioning of the Auston switch outside of the vacuum chamber of the intensifier and laser-produced plasma. The use of stripline for the pulse generator enabled easy access to the line conductor for biasing both the switch and the gated intensifier.

The Auston switch pulse generator produced a fast rising edge. The falling edge, to turn the device off, was produced by forming a transmission line loop (Figure 1.24). The positive and negative 'edges' emanate from the switch and run through each other on the opposite side of the loop. The photocathode is positioned so that the negative edge arrives some 100 ps before the positive. The DC bias to the switch is removed from the loop with a single chip capacitor. The full circuit is shown in Figure 1.24. In addition, the photocathode is biased to about +100 volts. The cathode response to various bias voltages is shown in Figure 1.25. The 100 V bias ensures that photoelectrons produced when the gate is off cannot drift into the image tube.

The gated intensifier is a proximity-focused device. Consequently it is subject to loss of resolution due to the lateral thermal spread of electrons as they pass from element to element. Photo-electrons produced by Caesium Iodide have a spread, V_a , of 1.6 eV. The corresponding thermal spread distance at the mesh, α , is related to the photocathode to mesh distance, s , by

$$s = \frac{\alpha}{2} \left(\frac{V_a}{V_{th}} \right)^{\frac{1}{2}}$$

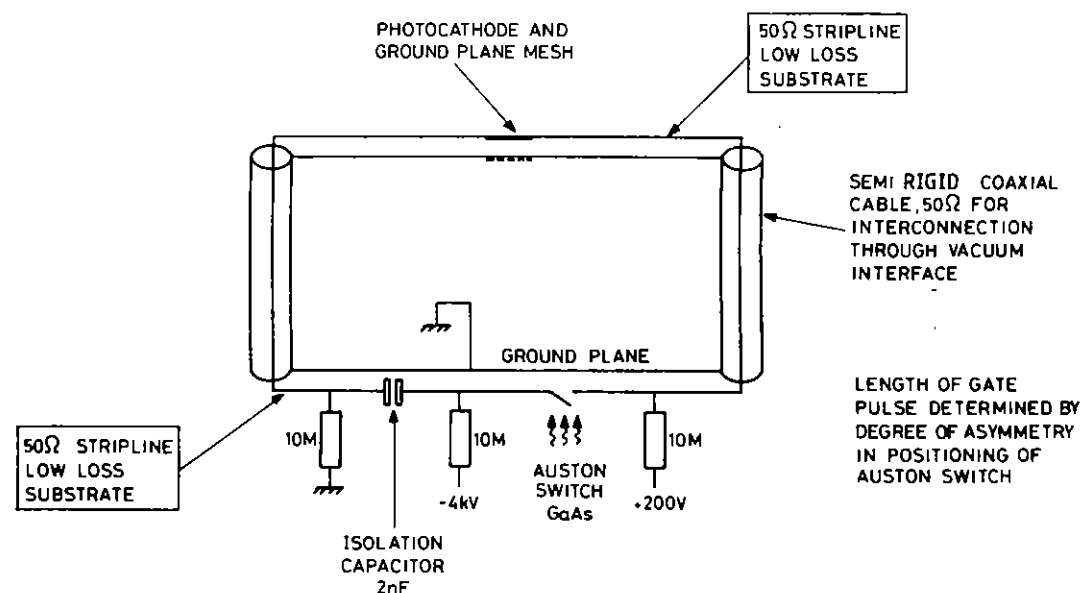


Fig. 1.24 Gate pulse generator and connection to the camera.

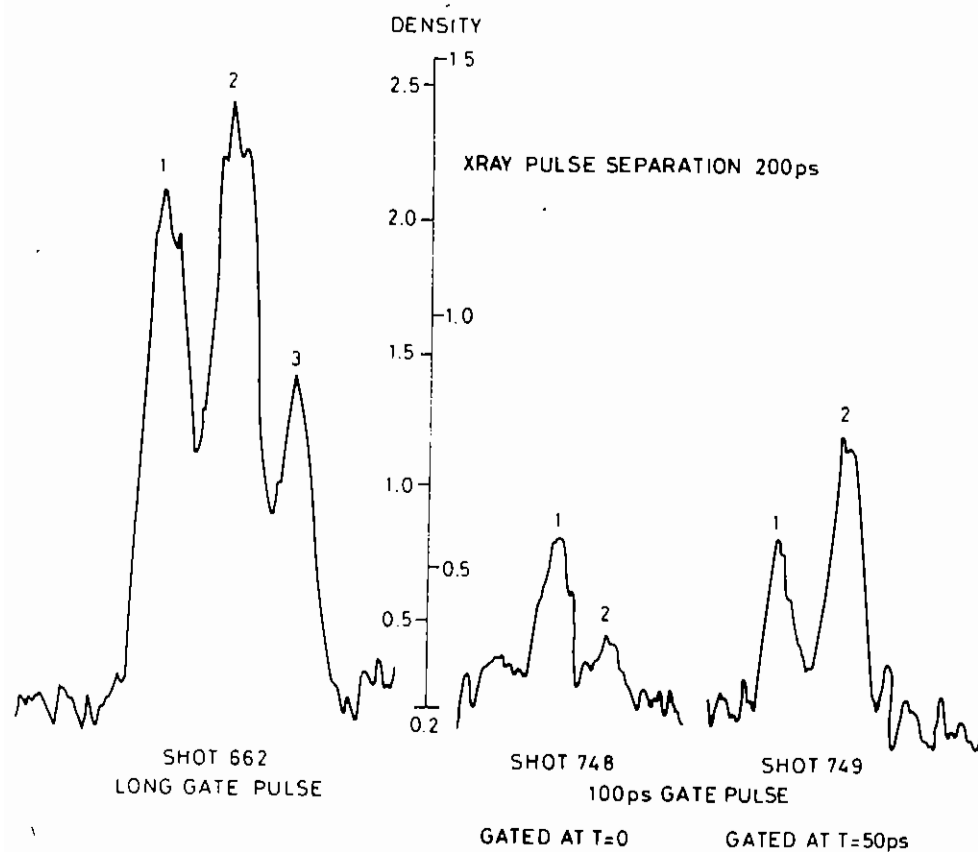
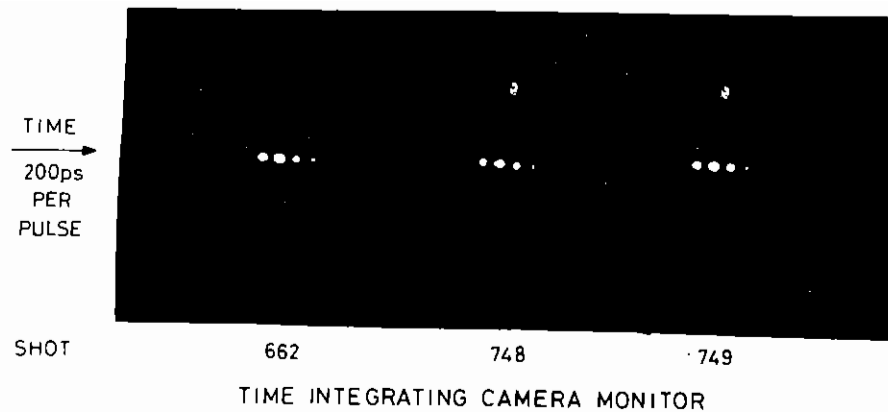


Fig. 1.25 Micro densitometer traces grating intensifier images.

where V_g is the gate voltage, if the transit time is less than the gate time.

The stripline characteristic impedance is 50Ω , for compatibility with coaxial cables. Consequently the width to separation ratio, w/s , of the stripline is ~ 5 . The number of resolvable picture elements across the photocathode is w/a . Hence

$$\frac{w}{a} = 2.5 \left(\frac{V_g}{V_{th}} \right)^{\frac{1}{2}}$$

With the 2 kV gate voltage available this gives 90 picture points.

The requirement that the electron transit time to the mesh be less than the gate time fixes

$$\left(\frac{2 s^2 m}{eV} \right)^{\frac{1}{2}} < t_{gate}$$

$$\text{Therefore } s < 0.94 \left(\frac{V_g}{1000} \right) \left(\frac{t}{100 \text{ ps}} \right) \text{ mm}$$

ie $s < 1.35 \text{ mm}$ at 2 kV.

In order to produce more picture elements it is desirable to increase the gate current. A reduction of stripline impedance is more beneficial than increasing the gate voltage as

$$\frac{w}{a} \approx \frac{125}{Z_0} \left(\frac{V_g}{V_{th}} \right)^{\frac{1}{2}}$$

In order to measure the gating time of the intensifier, three plasma X-ray sources were produced from one laser pulse. A 5 J, 100 ps, 1.05 μm laser was split up with an angled etalon to produce several focal spots on a solid copper target. Focusing with a 30 cm lens produced three bright X-

ray sources separated in space by 400 μm and in time by 200 ps.

Part of the same beam was extracted before the etalon and used to drive the Auston switch. A variable delay was placed in this beam. The reflectivity of the etalon elements was such that the relative intensity of the first four pulses was 0.3, 0.49, 0.15 and 0.044, although the fourth pulse was not visible. The plasma X-ray sources were imaged onto the photocathode with a 50 μm pinhole using 1:1 magnification. A separate ungated image intensifier using a microchannel plate for photon detection was used to monitor the generation of the X-ray sources.

The spectrum of Cu is a quasi-continuum with broad bands corresponding to unresolved L shell transition arrays. The duration of X-ray emission is much longer than 100 ps. The intensifier was filtered with 6 μm Al which transmits 40% in the region $\lambda > 9 \text{ \AA}$, and begins to transmit again at $\lambda < 4 \text{ \AA}$.

The intensifier can record images when gated off because the photocathode has a transmission of 30% to 8 $\text{\AA}$ X-rays which liberate photo-electrons from the microchannel plate, albeit with a much reduced quantum efficiency. Tests without gating the cathodes gave a null signal corresponding to a density < 0.3 on Figure 1.26. The contrast between a long gate pulse and no gate is therefore ~ 100 .

The X-ray emission from each source rises in ~ 50 ps (for a 100 ps FWHM laser pulse) but decays more slowly. Consequently the first pulse is still emitting when the second pulse begins. As a result, if the intensifier is gated on during the second pulse it will still see the first pulse. However, if the gate is turned off just before the fast rising edge of the second X-ray pulse, then only the first is seen. In this way the turn-off time may be estimated. In order to measure turn-on times with confidence, an X-ray source with fast turn-off is required.

Some results are shown in Figure 1.26. Three shots are shown, #662, #748, and #749. At the top of Figure 1.26 are the results from a separate time-integrating monitor. On shot #662 the gate pulse was made very long and

CsI photocathode illuminated with Ti characteristic X rays $\approx 10^8 \text{ photons } \text{str}^{-1} \text{ s}^{-1}$

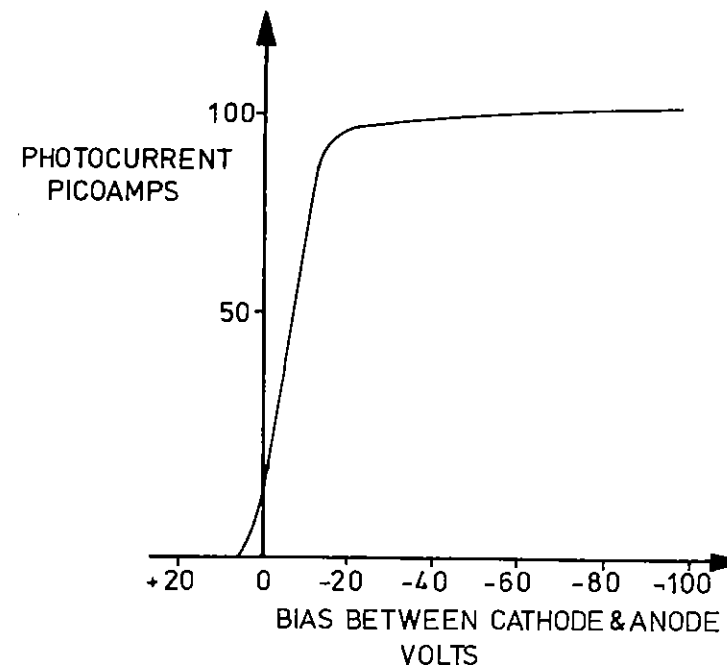


Fig. 1.26 Photocurrent as a function of extraction field.

the intensifier has imaged three sources. Shots #748 and #749 were gated for 100 ps, shot #749 being gated 50 ps later. It can be seen quite clearly that by gating just 50 ps later, the intensity of the second source has risen considerably. In addition, the intensity of the signal with a short gate pulse is ~ 0.1 of that with a long gate pulse, indicating ~ 50 ps gating time. The turn-on time should be similar as the circuit is symmetric.

It is a pleasure to acknowledge the assistance of the MERLIN operating crew at AWRE in testing this device.

A K L Dymoke-Bradshaw, J D Kilkenny (Imperial College) and S Smith (AWRE Aldermaston)

1.7.3 Compact X-ray Streak Camera

The main reason for designing a miniaturised streak camera is the ever increasing competition for solid angle with the addition of new diagnostics in the multibeam target chamber. The greatest saving in space has been achieved by developing a 50 Ω line-driving sweep circuit thus enabling the electronics to be remote from the sweep plates. The avalanche transistor circuit is able to drive a 1.5 kV pulse with 1 ns rise time into 50 Ω cable. Two such circuits, delivering positive and negative going pulses form the heart of the sweep circuit. The maximum sweep rate is ~ 20 ps/mm however tests on an optical streak camera at Imperial College have demonstrated that a sweep rate of ~ 5 ps/mm is possible although this is of little use for the X-ray streak camera.

The solid angle occupied by the new streak camera is ~ 0.2 sr to be compared with ~ 0.7 sr for the old TAL streak camera, with both photocathodes at a nominal 20 cm from the target. The new streak camera is sufficiently small to be re-entrant into a 250 mm port and hence can go arbitrarily close to the plasma.

Other improvements are:

(i) Modified electron optics giving reduced magnifications, thus allowing a larger area of the photocathode to be used.

(ii) Provision of a new intensifier with automatic fiducial marking and compatibility with 612 polaroid film.

(iii) Provision of timing and sweep monitor output signals to facilitate easy timing.

J Hares (Imperial College)

1.7.4 X-ray Pinhole Camera with Image Intensifier Recording

Early experiments at the Queen's University, Belfast (1.10) demonstrated the feasibility of using an aluminised phosphor screen as an optical transducer. High efficiency conversion of moderate energy X-rays (~ 1 keV) into visible light photons was obtained in a P11 phosphor (areal density ~ 5 mgs cm^{-2}). This light output could then be amplified in a conventional image intensifier, the complete system providing a simple and efficient X-ray recording medium.

Using this technique a high magnification X-ray pinhole camera is under construction at the CLF for use in the forthcoming high density ablative implosion experiments. The principal features of the design are:

(i) Image recording at the target chamber wall using the phosphor/intensifier scheme outlined above.

(ii) Variable magnification from $\times 20$ - $\times 50$.

(iii) Pinhole alignment achieved by external micrometer drives acting on the pinhole support tube held in a flexible steel diaphragm.

(iv) The large size of the detection phosphor screen (5 cm) facilitates easy alignment even at high magnification (900 μm field of view at $\times 5$ magnification).

Using this instrument and the technique of short pulse X-ray backlighting (1.02) we hope to record good quality, high magnification radiographs of the implosion experiments.

A J Cole, E Madraszek (RAL) and C L S Lewis (QUB)

1.7.5 Inertial Target Mounting System for Microballoons

A beneficial international collaboration with the Max Planck Institut fur Quantenoptik in West Germany has resulted in a novel prototype target release device developed at MPIQ being adapted for use in trials on our six beam spherical target irradiation facility.

Microballoon targets have generally been mounted on 5 to 10 μm glass or carbon fibre stalks using epoxy resin. The stalk and its glue joint introduce a significant non spherical perturbation when the target is imploded and this is thought to degrade the compression. In the new 'mousetrap' system developed at MPIQ (see Figure 1.27) the target support is a flexible metal arm with a small graphite rod on which the microballoon rests without any glue bond. The mounting arm has a motorized X-Y-Z positioning mechanism. There is also a spring loaded striker arm.

The microballoon target is injected from a syringe through a stainless steel capillary tube to a graphite nozzle positioned directly over the target pad. The target is then aligned to a point 5 μm above the focal point of the lenses. The system is fired by releasing the striker arm with a rotary solenoid. After release the striker arm passes a magnetic sensor which initiates the laser firing sequence. The timing is set up so that the laser pulse hits the target 1 ms after the striker arm has hit the target support arm. The combined speed of the support arm and the striker arm is $\sim 20 \text{ ms}^{-1}$ so the support arm is $\sim 2 \text{ cm}$ below the target and the target has dropped 5 μm under gravity when the laser pulse arrives at the target.

The system with adaptation to the CLF six beam target chamber has been tested in preparation for a trial run in experiments scheduled for May 1983.

Figure 1.28(a) shows a microballoon of about 200 μm diameter sitting on the pad and Figure 1.28(b) shows the microballoon's position 9.25 ms after removal of the support pad.

A R Damerell, C Brown, R W W Wyatt (RAL), R Sigel, G Tsakaris and W Haas (Max Planck Institut fur Quantenoptik, W Germany)

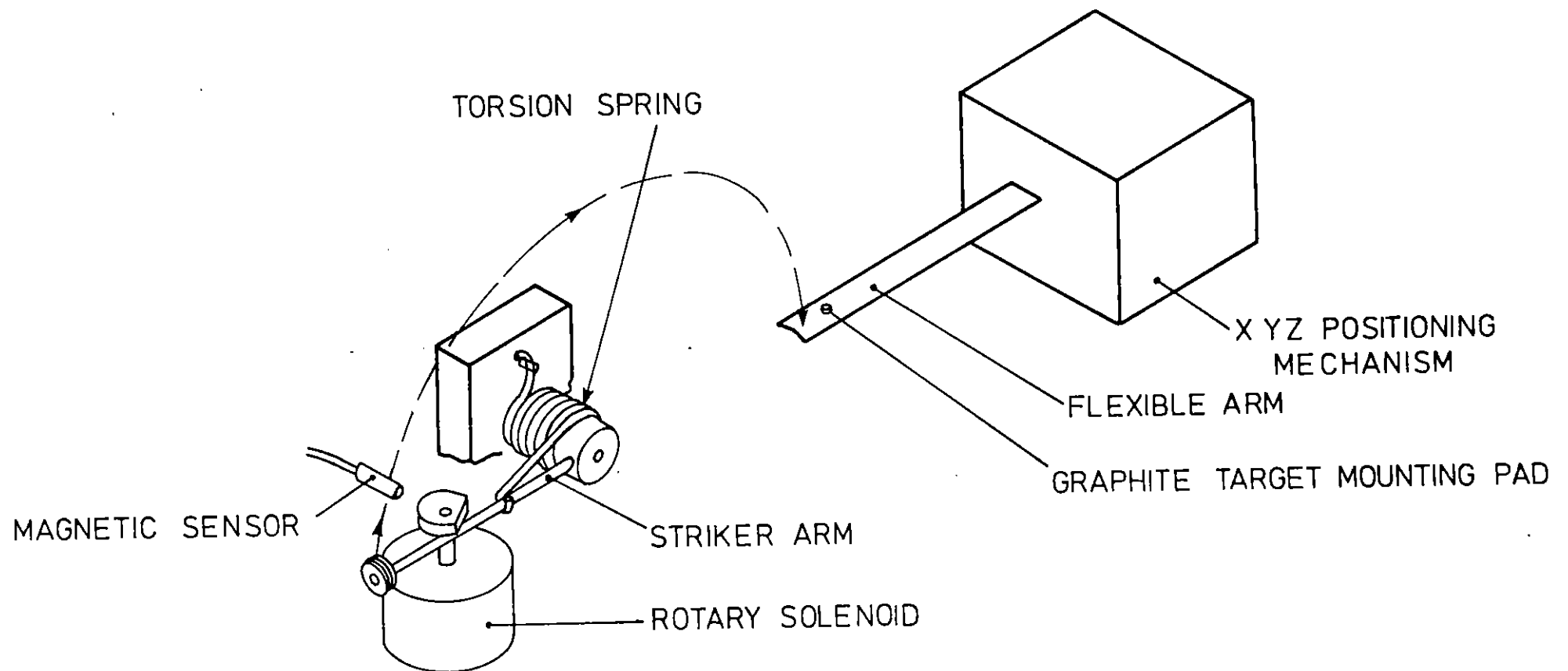


Fig. 1.27 Diagrammatic representation of the "Inertial Target Mount".

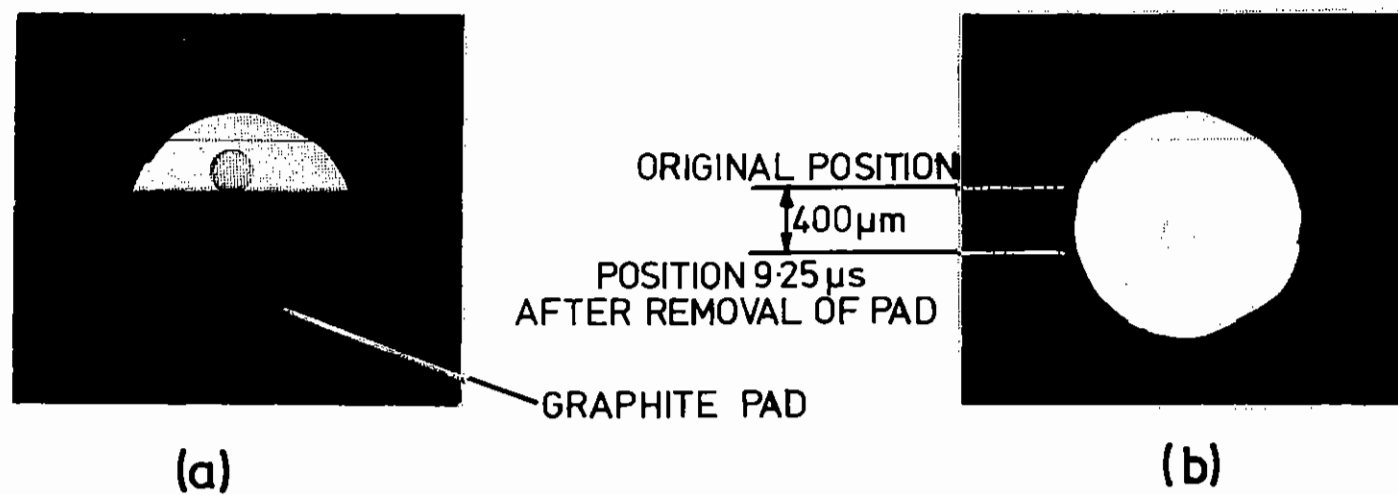


Fig. 1.28 Polaroid photographs showing position of target

- a) Before removal of graphite pad.
- b) 9.25 μs after removal of pad.

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CHAPTER 2 GAS LASER DEVELOPMENT

INDEX

2.1	INTRODUCTION	page 2.1
2.2	THE SPRITE KrF LASER	page 2.2
	2.2.1	System Development
	2.2.2	Long Pulse Laser Performance
	2.2.3	Sprite Pulse Compression Project
2.3	KrF LASER DEVELOPMENT	page 2.18
	2.3.1	Large Aperture, Long Pulse KrF Discharge Laser Feasibility Study
	2.3.2	UV Laser Damage Studies
	2.3.3	UV Phase Conjugation
	2.3.4	Pulse Compression and Phase Conjugation by SBS
REFERENCES		page 2.40

CHAPTER EDITOR: F O'Neill

CHAPTER 2 GAS LASER DEVELOPMENT

2.1 Introduction

The main advance in gas laser development during the 1982 reporting year has been the successful completion of the Sprite laser project. Laser oscillation was first achieved in May 1982 and by the end of the year an output pulse of ~ 220 J with a duration ~ 45 ns had been demonstrated. Low divergence (< 100 μ R) beams of > 160 J had also been obtained using 25 cm diameter unstable resonator optics on the laser.

With the successful operation of Sprite our efforts have now switched to solving the problem of pulse compression (and associated power multiplication) of KrF lasers. In particular we have started work on a project the aim of which is to compress the output of the Sprite laser into a 1 ns pulse with an efficiency of $> 50\%$, ie > 100 J pulse energy. This will be achieved using a scheme of staged optical multiplexing of KrF and methane Raman amplifiers and preliminary experiments on the Raman devices have just begun. According to our schedule the 268 nm pulse compressed Sprite laser will be available for target experiments in 2-3 years time. This UV target facility will act as a test bed for a 4-5 kJ, 0.6 ns multibeam 268 nm compression laser, OBERON, which will be constructed if the technology of the Sprite system proves successful.

Research work on phase conjugation has continued during the reporting year. Noteworthy here has been the shortest wavelength demonstration yet of phase conjugation using a 193 nm ArF laser. A narrow line (~ 0.5 cm^{-1}) ArF discharge laser had to be specially developed to demonstrate phase conjugate reflection from trifluoropropanol and hexafluoropropanol SBS mirrors.

F O'Neill

2.2 The Sprite KrF Laser

2.2.1 System Development

Over the past year the main improvements to the Sprite KrF laser system have been the installation of laser triggering on the PFL gaps and the development of a reliable foil glueing technique for the main pressure foil. These improvements have resulted in a large increase in the reliability of the machine which has to date been fired several tens of times at full power and several hundred times at reduced power.

2.2.1(a) The Laser Triggered PFL Gaps

Previous LFC annual reports have discussed the operation of the Sprite machine with the PFL gaps firing in self-break mode. Figure 2.01 shows the block diagram of Sprite with laser triggering. The laser used for triggering is a small KrF discharge laser which was constructed in-house using the internal peaking capacitor design. This laser gives a maximum output energy of 100 mJ in a 20 ns FWHM pulse with a 2 ns (10-90%) risetime. An $m = 20$ positive branch confocal resonator is used which gives an output beam divergence of better than 0.5 mRad. The laser is switched by a midplane-triggered spark gap which is triggered from a 50 kV, 2 ns risetime pulse derived from a capacitive divider measuring the voltage on the output terminal of the main spark gap in the water capacitor. The inductors charging the four Sprite PFL's and the delay line to the laser are chosen so that the laser pulse arrives at the output gaps when the voltage on the line is a maximum. The jitter in the laser output pulse (relative to the trigger pulse) was measured to be ± 2 ns (jitter is here defined as the standard deviation of the delay distribution).

The laser beam is split into two using a 50% beam splitter and into four using a further two beam splitters and steered into the four PFL gaps using AR coated 90° prisms.

The laser triggering beam enters the PFL through a fused silica window as shown in Figure 2.02 through 150 mm of water before entering the SF_6

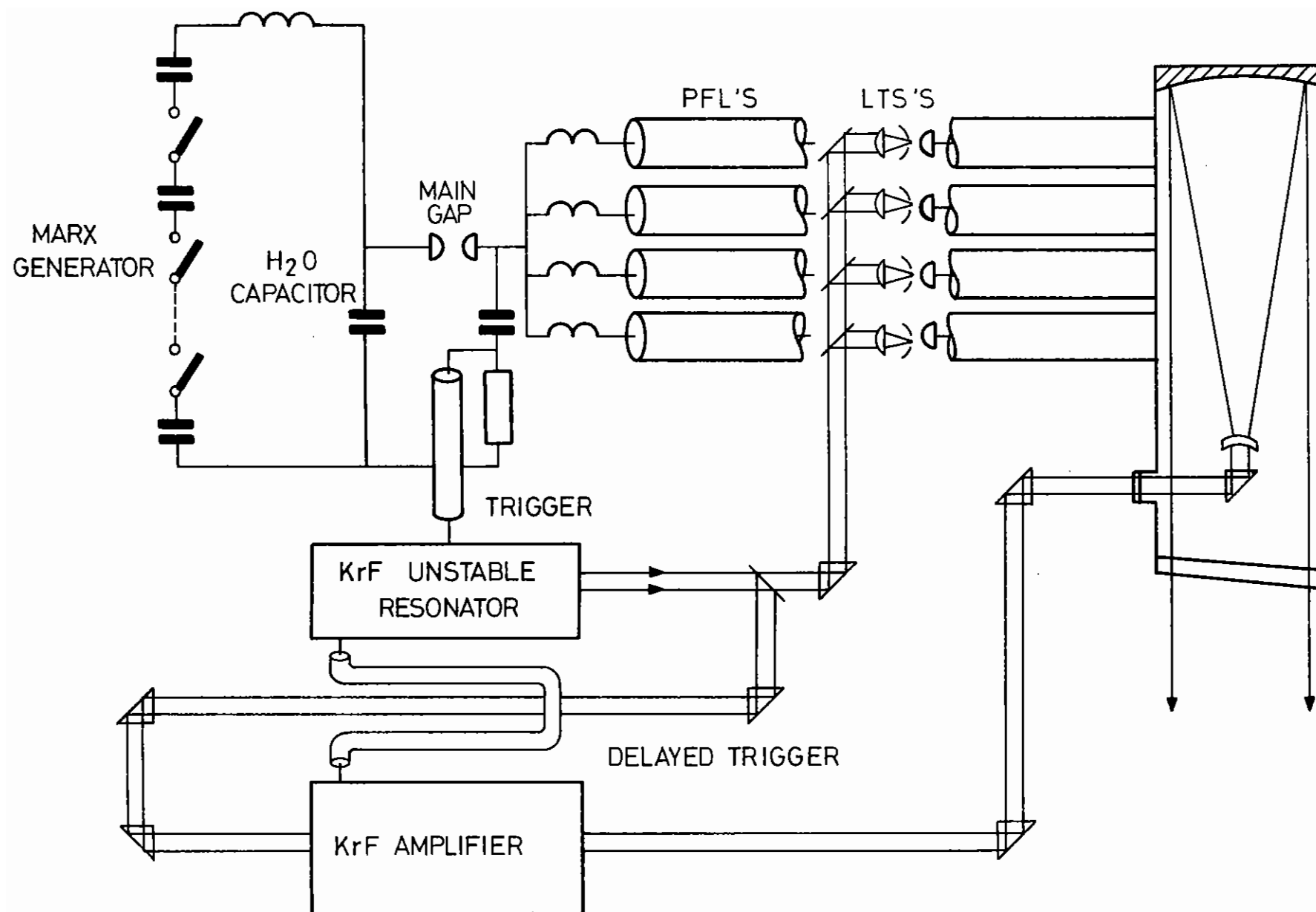


Fig. 2.01 The Sprite laser triggering system and optical layout for injection locking.

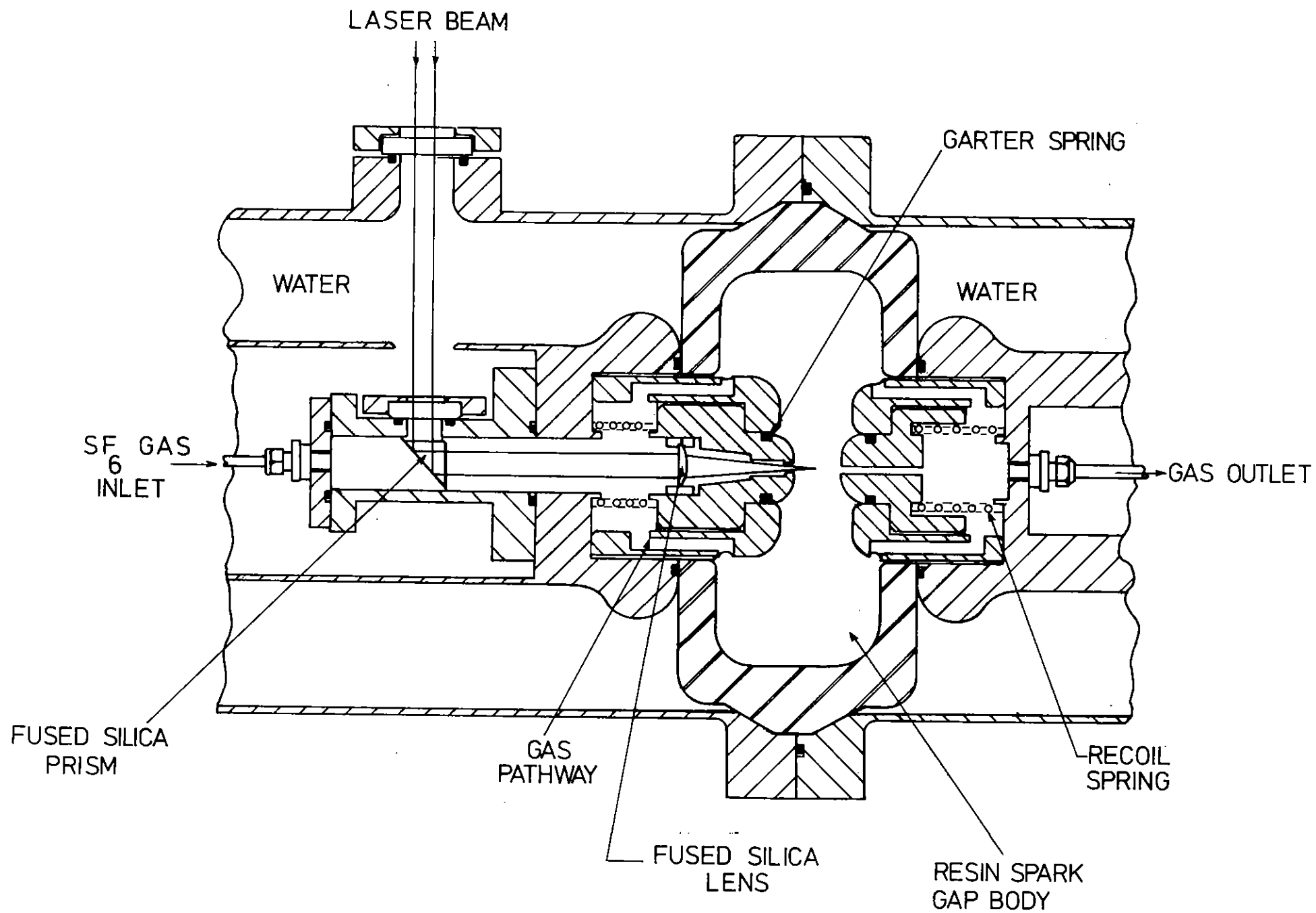


Fig. 2.02 The laser triggered switch. The left hand side is the charged side.

chamber through another fused silica window. Fortunately the absorption of 248 nm radiation in the filtered, de-ionized water is rather small. Alignment is initially achieved using a HeNe laser beam which is collinear with the KrF beam. This beam is arranged to enter the PFL window normally and when correctly aligned the HeNe beam passes through the anode hole and can be seen lighting up the plastic gas outlet tube of the spark gap. Unfortunately it is not possible to see into the gaps when they are assembled to ascertain whether the gas is actually breaking down or not, however this can be done acoustically by breaking into the gas outlet tube (with the SF₆ pressure at 1 bar) and listening for the characteristic "crack" of laser breakdown. The laser intensity at the focus is given by $I = 4E/\pi t f^2 \theta^2$ where E is the energy per pulse, t the pulse duration, f is the focal length of the lens and θ the full angle beam divergence. Setting I equal to the breakdown threshold in SF₆ of approximately 2 GW cm⁻² (2.01) we calculate the minimum laser energy required for breakdown as 400 μJ from the 20 ns pulse length discharge laser.

The main result of laser triggering was to give a dramatic improvement in the timing of the PFL gap firing over the self-break performance. This is clearly shown in Figure 2.03 which shows the charging waveforms of the PFL's with and without laser triggering. Quantitative measurements of delay and jitter were made by triggering two 7844 dual beam oscilloscopes from the laser beam. Each channel measured a PFL voltage as recorded by a capacitance voltage monitor situated about 20 cm from the output gap. The rapid change in the sign of dV/dt as the V/2 step propagates past the voltage monitor gives a sufficiently easily determined point on the waveform to measure the switch closure time to an accuracy of about ± 0.5 ns. The absolute t = 0 delay time is positioned on the oscilloscope traces by placing a photodiode at the PFL laser entrance window and recording the pulse arrival time on the voltage monitor trace using the same cable length as that used by the voltage monitor. Allowance is then made for the propagation time of the laser pulse to the gap centre and the propagation time of the EM pulse from the gap in the water dielectric to the voltage monitor point. The absolute accuracy of the delay time measurement is thought to be about ± 1.0 ns. Figure 2.04 shows the variation of delay with gap pressure which, using the self breakdown

characteristic for the gap, may be related to the fraction of self-breakdown voltage at which the gap triggers. The bars in Figure 2.04 indicate ± 1 σ in the distribution taken over typically 4 shots which gives 16 measurements of delay. Figure 2.05 gives the variation of delay and jitter with laser energy per gap. The threshold agrees well with the calculated value of 400 μJ. The relatively slow variation in delay and jitter with laser energy is an extremely useful property of the triggering system and means that it is relatively insensitive to slight misalignment, dust on optics etc. Our experience over several hundred shots with this system is that it is rugged and extremely reliable and is an important technological advance for high power laser system.

D Baker, D Craddock, C B Edwards, F O'Neill and M J Shaw (RAL)

2.2.1(b) E-Beam Foil Bonding

A considerable amount of effort in the past year has gone into finding the most suitable bonding technique for the main pressure foils on the laser cell. At the outset it was decided to use a bonding technique for the pressure foils for Sprite since previous experience had shown that with the use of pre-foils a pressure foil could last for thousands of shots and thus frequent replacement was not required. Initially Inconel foil 25 μm thick was bonded to the stainless laser cell using epoxy adhesive. Experience with this foil was not good since although the main bond did not fail, this thickness of Inconel was found to be prone to develop small leaks especially near to support structure ribs where the foil was sharply creased. A second cell was tried using untreated 25 μm titanium foil which was bonded to the cell with cyanoacrylate adhesive (Eastman 910). Although this foil was superior with regard to pin-hole formation, the bonding technique was unsuccessful due to the poor shock resistance of cyanoacrylate bonds. The foils experience a pulsed loading when the e-beams are fired into the laser gas. The final and successful foil and bonding technique used employed a 35 μm titanium foil which was pre-treated in a hydrogen peroxide bath (2.02) prior to bonding to produce a uniform oxide layer on the surface of the titanium which had far superior bonding properties to that of the bare metal. An added advantage of this

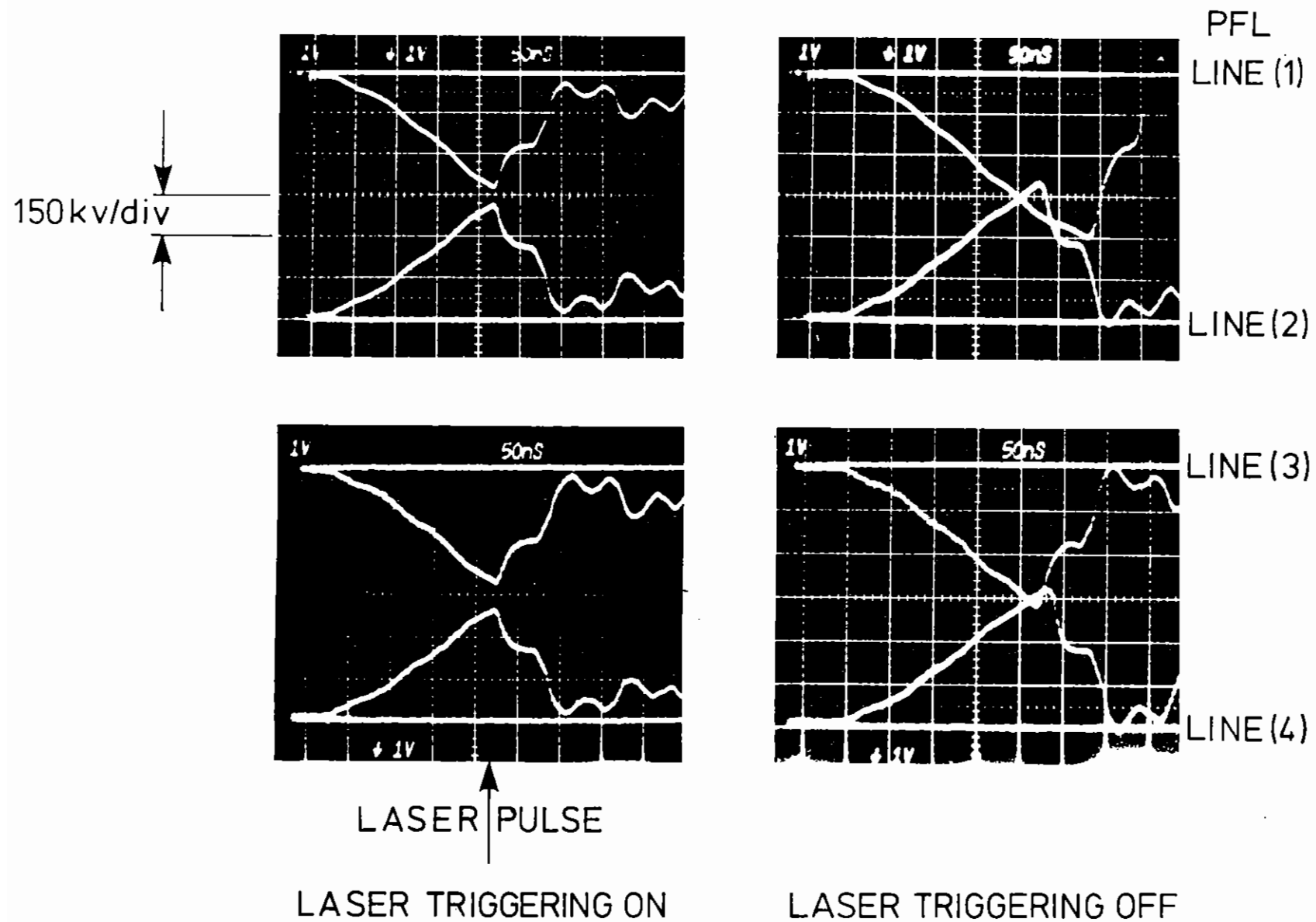


Fig. 2.03 Voltage waveforms of the four PFL's as measured by the capacitance voltage monitor with and without laser triggering.

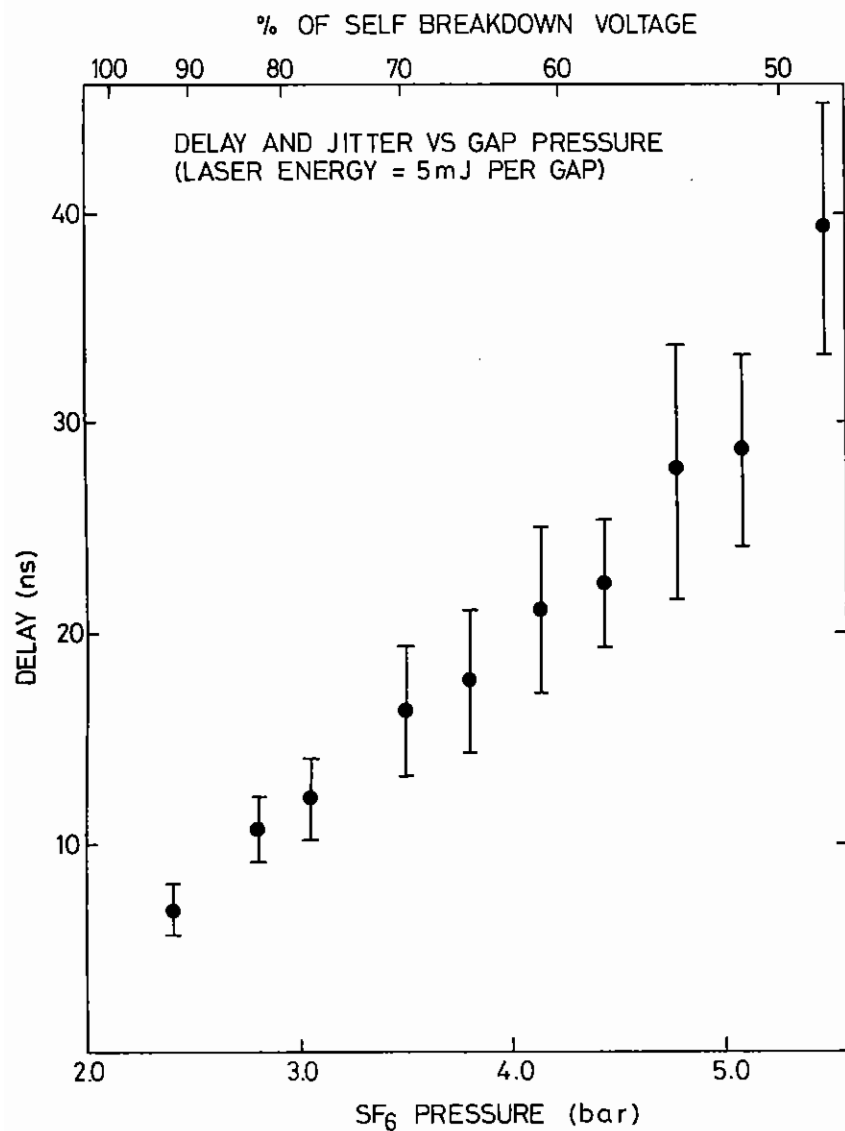


Fig. 2.04 Delay between laser pulse entering gap and onset of voltage collapse as a function of gap pressure. Error bars indicate jitter (± 1 standard deviation of the delay distribution).

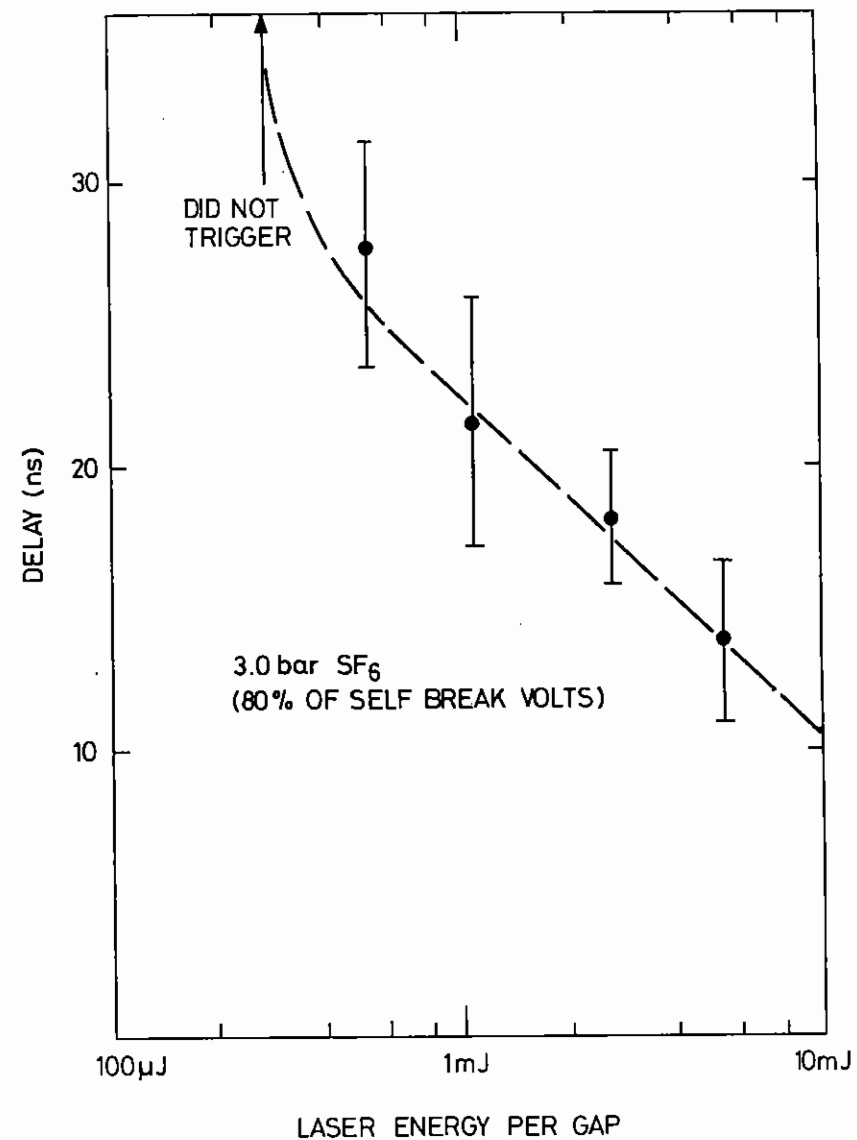


Fig. 2.05 Delay and jitter of the laser triggered switch as a function of laser energy at constant gas pressure.

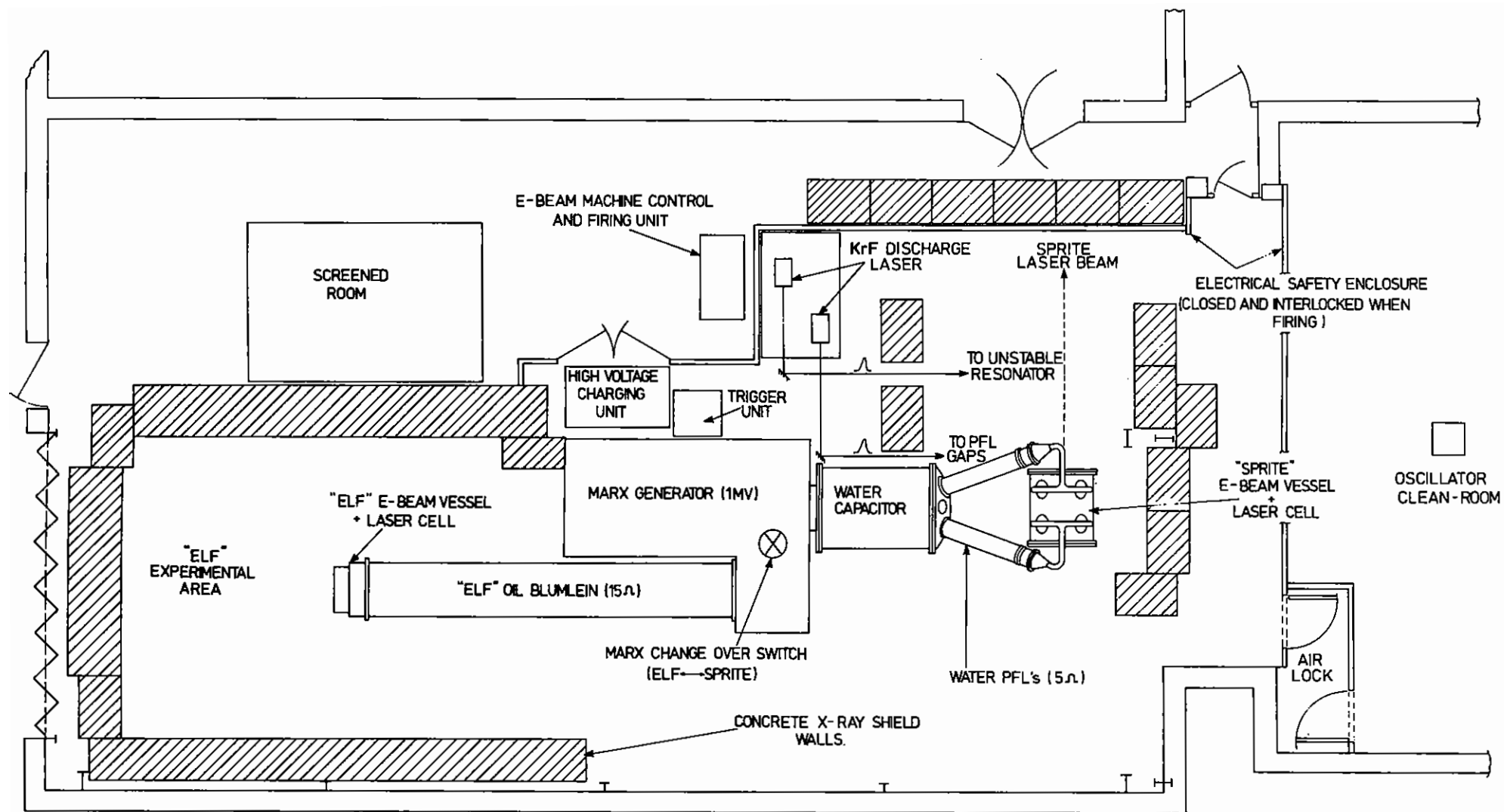


Fig. 2.06 Layout of Sprite laser laboratory.

pre-treatment was that the foils were rendered dull, the colour varying between light straw to dark grey depending on thickness of the oxide layer, with a matt surface. Thus the foil reflectivity is expected to be less than for the bare metal foils. The adhesive employed is Redux 410 (Ciba-Geigy) which remains slightly flexible on setting and appears to have excellent shock resistance. The whole pre-treatment and bonding procedure is done in-house and has, after some practice, been found to be relatively easy. A detailed description of the procedure is available from the Chemical Technology Group (2.03) at Rutherford Appleton Lab. One such cell foiled in this way has by now withstood many hundreds of shots without failure.

D Baker, D Craddock, J Landgridge and M Shaw (RAL)

2.2.1(c) E-Beam Performance

The Sprite electron beam system has so far met the specification figures and with present modifications may well exceed them. The single most important item determining the e-beam performance is the laser triggering system. This allows routine charging of the PFL's to over 900 kV giving 10 kJ of stored energy in the lines. The Sprite system is only properly matched when all four lines fire simultaneously and thus laser triggering insures proper matched performance of the system. Prior to laser triggering, operation at high voltages almost always led to breakdowns due to reflected energy from the mismatched diodes.

The diode inductance gives the expected 10-90% current rise-time of 15 ns with a 40 ns flat top region. Integration of the I and V traces gives a diode energy of 7.5 kJ (with lines storing 10 kJ). Of this 7.5 kJ only about 3 kJ gets into the gas. This is somewhat lower than expected from trial experiments however it is sufficient to produce a laser output energy in excess of 200 J. A possible explanation of the 1 ns deposition efficiency has been found from examining the pre-foils after a hundred shots or so. Marking of the foil opposite the edges of the cathode holder showed that appreciable emission had been occurring from the cathode holder itself rather than the cathode emitter plate. This emission at large

angles to the main foil probably would not get into the laser gas. The edge of the cathode holder has been modified to lower the field there and the cathode spikes have been taken back from the edge of the holder. Furthermore the prefoil has been changed from 12 μ m Inconel to 12 μ m titanium which should lower losses and improve electron trajectories. It is hoped that these modifications will lead to increased deposition efficiency for the system.

Recently the system has been stripped down for inspection after several hundred shots. Apart from the cathode holder problem mentioned above there was no evidence of any major breakdown problems in the system. In particular the laser triggered switches looked very clean with no evidence of surface tracking which had previously been a problem when running the gaps under self-break conditions.

C B Edwards, F O'Neill and M J Shaw (RAL)

2.2.2 Long Pulse Laser Performance

Sprite first became operational as a laser in May 1982. Initial tests for lasing were performed using an optical resonator consisting of an aluminised rear reflector (back coated Al on fused quartz) and an uncoated quartz window aligned for use as an output coupler ($R \sim 10\%$). For these early experiments laser triggering was not installed on the PFL output gaps with the result that these gaps operated in the self-break mode and exhibited comparatively poor synchronisation (1σ jitter ~ 10 ns). Laser oscillation was however easily achieved on these runs and KrF output pulses of energy ≈ 100 J and duration ~ 40 ns were observed. During this period of operation the e-beams were pulsed to voltages in the range 450 - 500 kV and a survey of the X-ray emission from the diodes was carried out. Based on this survey we have designed a layout of concrete shielding blocks around Sprite as shown in Figure 2.06. This arrangement of blocks provides for shielding of operating personnel from X-rays and at the same time gives good access to the machine for diagnostic and maintenance purposes.

Following initial experimental work, some improvements were made to the Sprite hardware. In particular laser triggering was installed on the PFL output gaps and an improved method was developed for glueing the laser pressure foils to the stainless steel laser cell. As a result of these two changes the reliability of the machine was greatly improved. The laser triggered PFL gaps provided well synchronised (1 - 2 ns jitter) voltage pulses to the e-beam diodes resulting in good impedance matching and alleviating high voltage breakdown problems which previously occurred due to voltage pulse reflections. The improved method of foil glueing means that laser gas pressures up to 1500 torr can now be safely used and the glue joints can easily withstand the pressure pulse which occurs in the gas due to heating by the deposited e-beam energy. The laser system is now capable of operating for some hundreds of shots without maintenance and KrF output energies of > 200 J in pulse durations ~ 45 ns have been obtained for a deposited e-beam energy ~ 2.7 kJ. It is worth noting that the laser cell does not seriously contaminate the laser gas. The major materials in contact with the gas are stainless steel, aluminium and titanium. After proper passivation we find that the laser gas fills can be used for tens of shots without changing. The gas does however have to be circulated through a 1 μ m pore filter to remove particulate matter generated by each e-beam shot.

Eventually it is planned to use Sprite as the main power amplifier in a high energy (100 - 150 J) short pulse (~ 1 ns) UV laser system. This short pulse system, which is described in more detail in Section 2.2.3, will utilize angular optical multiplexing and Raman amplification to compress the energy available from the 60 ns Sprite system to a 1 ns pulse suitable for laser target experiments. Angular multiplexing of a KrF laser has previously been demonstrated by the group at Lawrence Livermore National Laboratory (2.04) and the application of this technique to Sprite should be straightforward. The same group has also studied the use of CH₄ Raman amplifiers with KrF pumping (2.05) but they have emphasised the backward Raman system and have experienced difficulties with this device at high intensities. We plan to optically multiplex CH₄ Raman amplifiers used in the forward mode running at lower intensities and there is little information available in the literature for these devices in the UV. Our

near term plan is therefore to use Sprite as a high power pump source for a CH₄ forward amplifier to carry out studies of gain and extraction efficiency.

In order to use the output of Sprite as a pump for a Raman amplifier it is necessary to produce a beam of narrow linewidth and low divergence. This will be done by using an unstable resonator which is injection locked to an oscillator having the desired beam characteristics. A further advantage of injection locking may be obtained if the injected power is sufficiently large to suppress ASE and start the desired output mode sooner than would otherwise occur, in which case greater output energy and fewer parasites may be obtained. In order to investigate the unstable resonator cavity and to demonstrate the injection locking property, the experimental system shown in Figure 2.01 was used. A small part of the laser beam which is used to laser trigger the PFL gaps is split off and delayed by 60 ns (the propagation time of the output lines from the PFL gaps to the diodes). This beam is then amplified by a single pass of a KrF discharge amplifier which is triggered from the oscillator thus ensuring accurate timing. After aperturing to a 10 mm diameter beam, 10 - 20 mJ with a divergence of better than 0.5 mRad are available for injecting into the Sprite unstable resonator. The time of injection may be varied by changing the optical path and the jitter is just that of the laser triggered PFL gaps viz at most a few ns. One problem is matching the mode of the injected beam to the Sprite beam which requires a 25 x magnification of the 10 mm diameter injected beam. This is done by using 4 passes of the m = 5 unstable resonator as shown in Figure 2.07. The injected beam enters the resonator through a diverging lens which also acts as the feedback mirror and has a 92% reflectivity coating. The focal length of the lens and the mirror are chosen to be identical (at 249 nm) and thus the injected beam undergoes a x 5 magnification in the first double pass to fill the feedback mirror and a further x 5 magnification to fill the cavity. The cavity fill time is 20 ns which is the same as the current rise time. The optimum timing for the injected pulse was found to be such that the injected pulse entered the cavity at the beginning of the current pulse thus filling the cavity at the time when the gain reached its maximum value. Figure 2.08 shows the effect of injection on the laser output pulse. Clearly

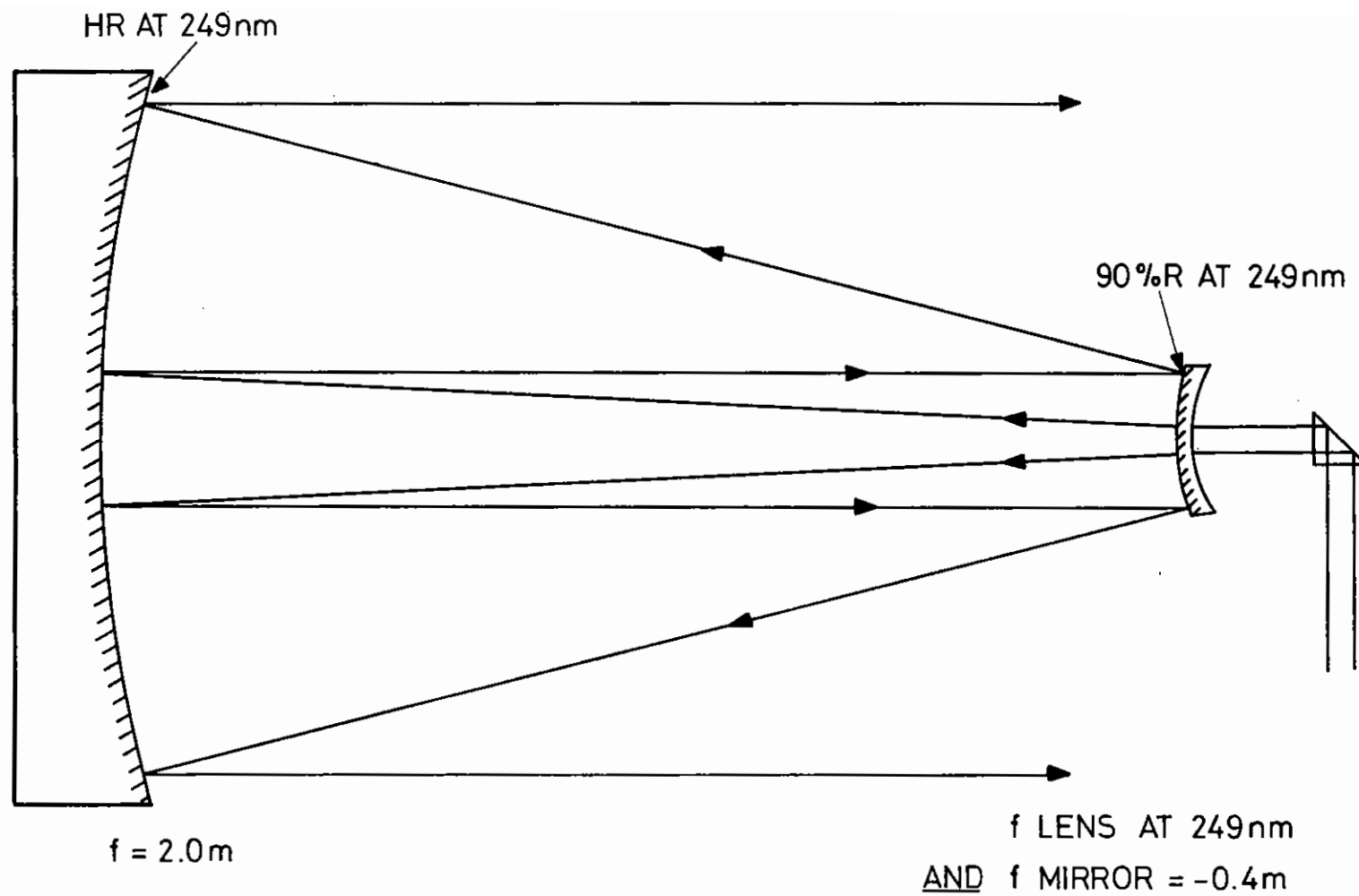


Fig. 2.07 Sprite M = 5 unstable resonator with four pass injection.

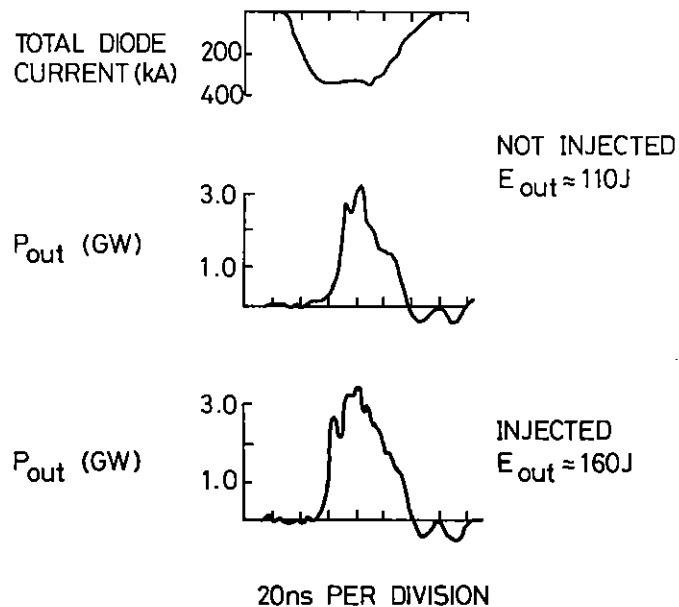


Fig. 2.08 Sprite laser output pulses with and without injection.

injection locking is achieving the desired effect with the injected case producing longer pulses with approximately 50% greater output energy. The divergence of the injected beam is reduced by the 25 x magnification and thus the output beam divergence should be $< 20 \mu R$ input. However in this system the output divergence is determined by the optical properties of the output window in transmission which limit the divergence to approximately $25 \mu R$. No attempt has been made to measure divergence, but Fresnel diffraction fringes in output burn patterns indicate that the divergence is much less than $100 \mu R$.

This experiment has demonstrated that a good beam quality injected laser can be used to initiate the Sprite pulse and produce low divergence high energy beams. It only remains for us to use a line narrowed ($\sim 0.2 \text{ cm}^{-1}$) injection laser, which has now been installed in our laboratory to produce high power Sprite beams suitable for pumping Raman amplifiers. This work is now in progress.

C B Edwards, F O'Neill and M J Shaw (RAL)

2.2.3 Sprite Pulse Compression Project

2.2.3(a) Project Overview

As has been described in the previous sections the Sprite KrF laser has produced $> 200 \text{ J}$ pulses of duration $\sim 45 \text{ ns}$. It is anticipated that with further development an energy output $\sim 300 \text{ J}$ in 60 ns pulses will be achieved. Even at this increased energy level, laser pulses of tens of nanoseconds duration are of only marginal interest for laser target irradiation experiments. Assuming that the 300 J , 60 ns laser pulse is focused to the smallest usable spot size of $100 \mu\text{m}$ diameter the irradiance on target will only be $6 \times 10^{13} \text{ W cm}^{-2}$. This is 1 - 2 orders of magnitude less intensity than will be required to generate uniform high density implosion effects in plasmas irradiated with $0.25 \mu\text{m}$ radiation (2.06, 2.07).

Obviously some method must be devised of increasing the intensity of KrF

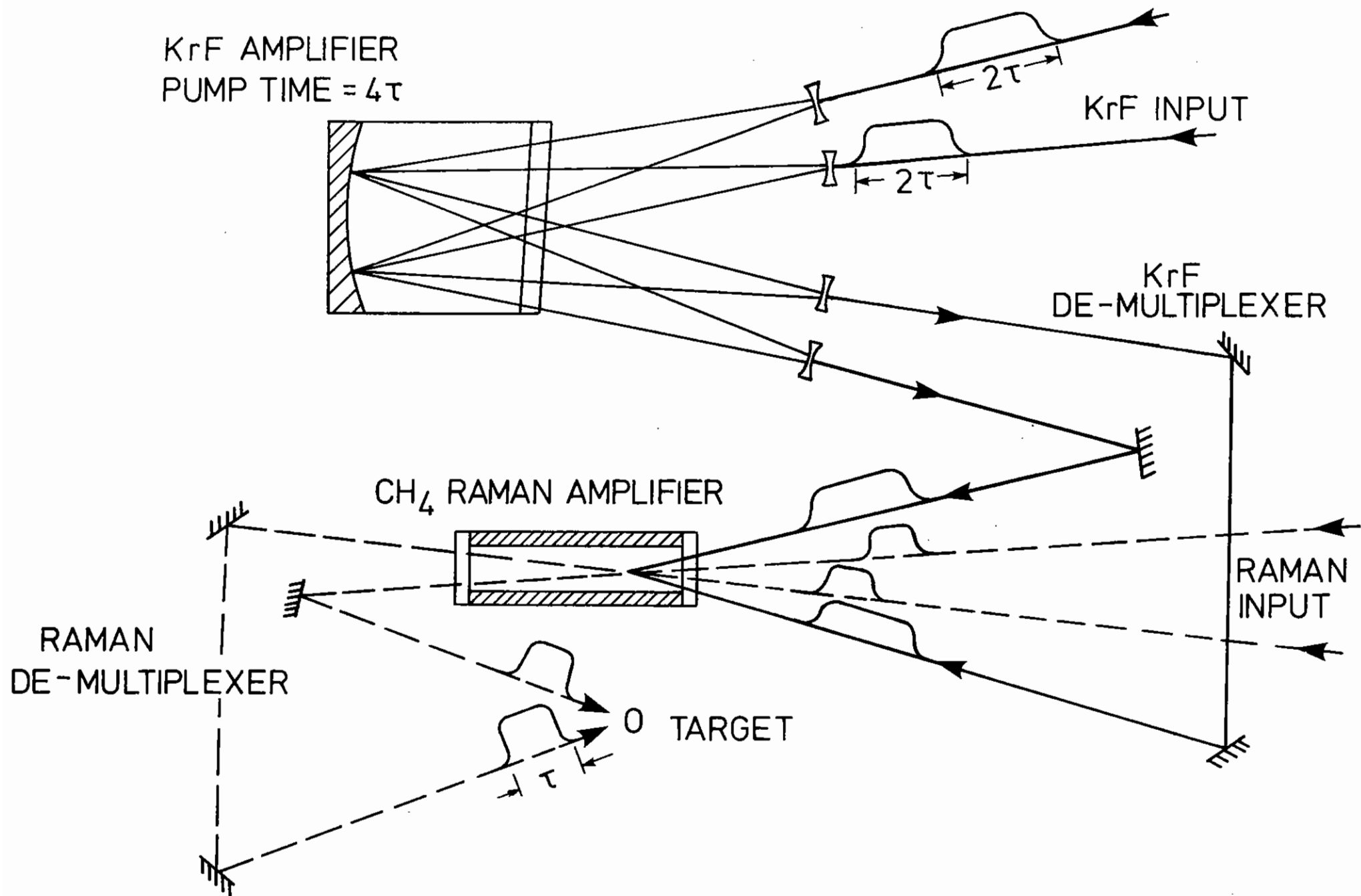
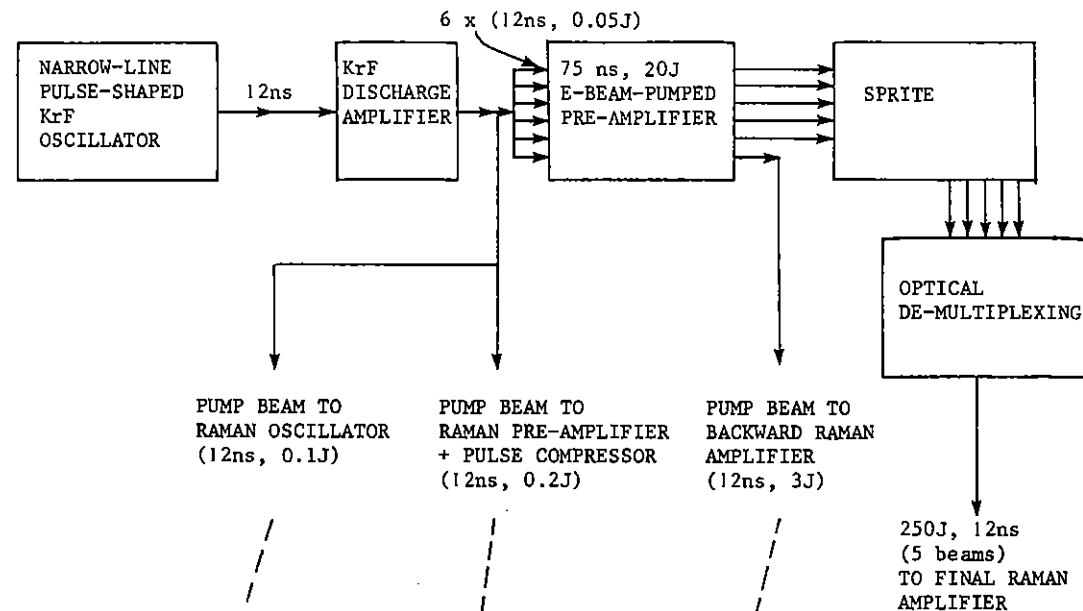


Fig. 2.09 Example of 2 x 2 staged multiplexer using KrF and a forward Raman amplifier. The full lines denote the KrF beams. The dashed lines show the 1st Stokes Raman beams.

(i) GENERATION OF 250J/12ns
PUMP PULSE FOR RAMAN AMPLIFIER



(ii) RAMAN AMPLIFIER CHAIN

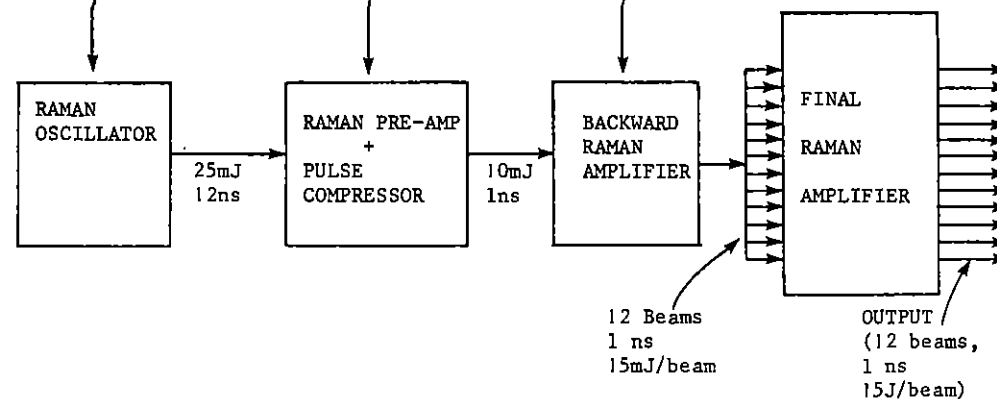


Fig. 2.10 Conceptual layout of 12 beam 268 nm 1ns laser using 5 x 12 staged multiplexing.

lasers before they can be usefully applied to target irradiation experiments. We plan to use a scheme of staged multiplexing to compress the output energy of Sprite from 60 ns to ~ 1 ns duration thus achieving the requisite order of magnitude increase in laser power. The pulse shortening scheme which we intend to use can be understood by referring to Figure 2.09 which shows the simplest case of a 2×2 staged multiplexer. The large aperture KrF amplifier is e-beam pumped for a time interval 4τ and is extracted sequentially by two angle encoded KrF laser pulses each of duration 2τ one delayed from the other by 2τ . Timing is adjusted so that the leading edge of the first extract pulse enters the amplifier at the start of the e-beam pumping pulse and the falling edge of the second extract pulse leaves the amplifier at the end of the pumping pulse. A double pass configuration is preferred for KrF because it optimises efficiency and reduces the possibility of parasitic oscillation (2.08). The extraction efficiency of the optically multiplexed KrF amplifier is determined by geometrical and gain considerations and can approach 100%. For example in non-optimised experiments at Lawrence Livermore National Laboratory > 20 J has been extracted from a multiplexed amplifier which produced ~ 25 J output when operated as a laser oscillator, ie the extraction efficiency was $\sim 80\%$ (2.04).

After amplification the two KrF pulses are decoded and appropriately delayed to simultaneously pump an optically pumped amplifier with a pump time of 2τ . This amplifier is itself optically multiplexed by two pulses of duration τ at the wavelength of the second amplifier. These new pulses are then decoded and directed on target in synchronism. The energy available from the large KrF amplifier has therefore been compressed from 4τ to τ with an intensity gain hopefully approaching 4.

The efficiency of the optically pumped amplifier is crucial for the success of this scheme as are its beam quality and pre-pulse level. Fortunately one system appears to possess all of these desirable characteristics viz the Raman amplifier. Gaseous H_2 Raman amplifiers have been reported (2.09, 2.10) to have conversion efficiencies $\sim 50\%$ to first Stokes and even higher when the pump beam was confined to a wave-guide (2.11). This latter technique is of considerable interest to multiplexed

Raman amplifiers because the effective beam overlap is greatly increased permitting near unity conversion even when the pump and Stokes beams travel at an angle to each other.

Using this principle we are at present designing a 5 x 12 staged multiplexing scheme to compress the expected 60 ns, 300 J output from Sprite into a 1 ns pulse of > 150 J energy at 268 nm using methane as the Raman medium. The general layout of the pulse compressor is shown in Figure 2.10 and basically consists of two parallel chains of laser amplifiers viz KrF and Raman (CH_4). The Raman amplifiers are pumped by multiplexed KrF beams and the final output from the system which will be used for target work will consist of twelve synchronised laser pulses each of 1 ns duration and ~ 15 J energy. The biggest single piece of hardware in the whole system is the Sprite KrF laser and this is now operational. The narrow line KrF oscillator and discharge amplifier have been designed to produce 0.5 J pulses with a line-width $\sim 0.2 \text{ cm}^{-1}$ and this equipment is now becoming operational. Work will start in the near future on constructing the 20 J e-beam-pumped preamplifier. This laser has been code-named GOBLIN and will be constructed using Sprite-type water-line technology and laser-triggered switching.

The Raman amplifier chain is the area where there is most uncertainty in the system design. The CH_4 oscillator and preamplifier have been built and are now being commissioned. In order to optimise the design of the final Raman amplifier we have constructed a light guided forward Raman amplifier which will be used for preliminary experiments. This amplifier will initially be pumped by the long pulse output of Sprite to measure the gain and saturation characteristics of CH_4 pumped at 249 nm.

More details concerning the progress on the KrF and Raman systems are given in the following sections.

2.2.3(b) Short Pulse Oscillator System

During the second half of the reporting year, part of the ELF laboratory, previously used as an experimental/workshop area, has been refurbished as

an air-conditioned clean room for use as a Krypton Fluoride/Raman oscillator laboratory. The area will be dedicated to the production of narrow linewidth, good beam quality KrF pulses as a front end for the Sprite multiplexer, and the generation of Raman shifted radiation at 268 nm as a short pulse (~ 1 ns) injection source for the multiplexed Raman preamplifier and final amplifier.

Considerable engineering effort has been directed to the construction of a high pressure (450 psi) methane, and fluorine handling system with associated controls and we now have a reliable installation, compatible with the stringent safety requirements, with considerable flexibility to allow for future developments and improvements.

The optical arrangement within the laboratory is based around a Lambda Physik type EMG250ES Krypton Fluoride oscillator/amplifier combination, a custom-built upgraded version of the EMG150. This laser produces pulses of 20 ns duration, with an energy of 0.5 Joules within a bandwidth of 0.3 cm^{-1} , with a beam divergence of better than $300 \mu\text{R}$ (full angle at half intensity). The repetition rate is up to 10 Hz. The device was delivered in February 1983 and is now installed in the oscillator room. At the present time, three Raman cells have been installed for use in the preliminary optical extraction efficiency experiments using methane at 268 nm, described in 2.2.3(c) below. The two high-pressure cells will be used for providing the necessary 268 nm extraction beams. The third cell is designed to operate at low pressure (< 50 psi) and contains an optical waveguide. The optical layout of the oscillator is shown in Figure 2.11.

For the immediate experiments, cell (1) will be used as a single pass superfluorescent 'oscillator' employing high small signal Raman gain ($\sim \exp 35$) achieved with a combination of high pressure (400 psi) and high intensity ($\sim 200 \text{ MW cm}^{-2}$). This requires a double telescoping configuration; firstly a cylindrical telescope to reduce the 3:1 aspect ratio of the laser output to square, followed by a conventional spherical telescope to give the required intensity. At such laser intensities it is not possible to employ coated optics; however since the optical pulse is polarised losses can be reduced using Brewster-angled windows on the Raman

cell. Despite the high gain for the first Stokes pulse, higher order processes are suppressed by the choice of plane wave geometry which prevents the phase matching of parametric processes which dominate high order Stokes production in a focused configuration.

As the 268 nm radiation passes out of cell (1), it is expanded in another telescope, and combined with 249 nm radiation from the KrF laser at a dichroic mirror. The beams then co-propagate through cell (2) which is operated at lower Raman gain. This cell is effectively a single pass forward amplifier, the 268 nm injection pulse extracting energy from the KrF pumping pulse to provide a suitable pulse for extraction of the waveguide Raman amplifier (cell (3)). At the lower intensities it is possible to employ coated optics which reduce reflection losses at beam steering optics to a minimum. Synchronisation of the pulses is achieved using variable optical delays as shown in Figure 2.11.

The spatial quality of the 268 nm radiation produced will be similar to that of the KrF discharge lasers. Whilst this will be adequate for the planned experiments in the near term, it will not be sufficient for the multiplexing experiments and possible target irradiation. It will also be noted that at this stage no pulse shortening is incorporated into the experimental configuration. It is proposed therefore either to construct a multipass guided Raman oscillator operated on a single transverse mode in conjunction with a preamplifier, or alternatively to use the second harmonic of a TEM₀₀ mode dye laser. The latter approach would be convenient for the generation of subnanosecond pulses, whilst the former would require Pockels cells to shorten the pulse. The 1 ns pulses would then be amplified efficiently in a Raman cell operated in the backward mode. The existing high pressure cell (2) would be suitable for this purpose.

2.2.3(c) The Light Guided Forward Raman Amplifier

The most important component in the Sprite pulse compression project is the Raman amplifier which allows for staging of the multiplexing and thus larger compressions using fewer mirrors. The properties of the overall

system will in the main be determined by the properties of this amplifier the most important of which are:

- (a) High conversion efficiency (target figure = 66%).
- (b) Good beam quality output from poor beam quality pump.
- (c) Large angular acceptance for both pump and output beams.
- (d) Low pre-pulse levels.
- (e) Freedom from parasitics and higher order processes.
- (f) Benign technology - eg low pressure, low intensities etc.

In order to achieve all of these properties it seems necessary to use light guided pump beams for the Raman amplifier. Previous work (2.11) using H₂ pumped at 1.06 μm has demonstrated $\approx 70\%$ conversion efficiency using this technique.

In order to test out this scheme at 0.25 μm using energies appropriate to the final system, the experiment shown in Figure 2.12 is being set up. This experiment is devised to test out all of the important properties of a light guided amplifier using 60 ns pump beams and a single 15 ns output beam. (The bandwidth of the proposed CH₄ Raman medium is ≈ 10 GHz, so extrapolation to 1 ns and even sub-nanosecond pulses should make no major difference in principle.) By using 60 ns pumping there is no need to multiplex Sprite, all that is necessary is to run the system as an injection-locked-oscillator (as we have already demonstrated) in order to produce a pump bandwidth < 10 GHz.

Multiple beam pumping is simulated by using a segmented plane mirror. This mirror comprises 56 individually adjustable plane rectangular mirrors 23 mm x 32 mm used at nominally 45° and each aligned to fill the entrance aperture of a 25 mm x 25 mm hollow light guide. The intensity at the entrance of the light guide is between 300 and 500 MW cm⁻². One advantage

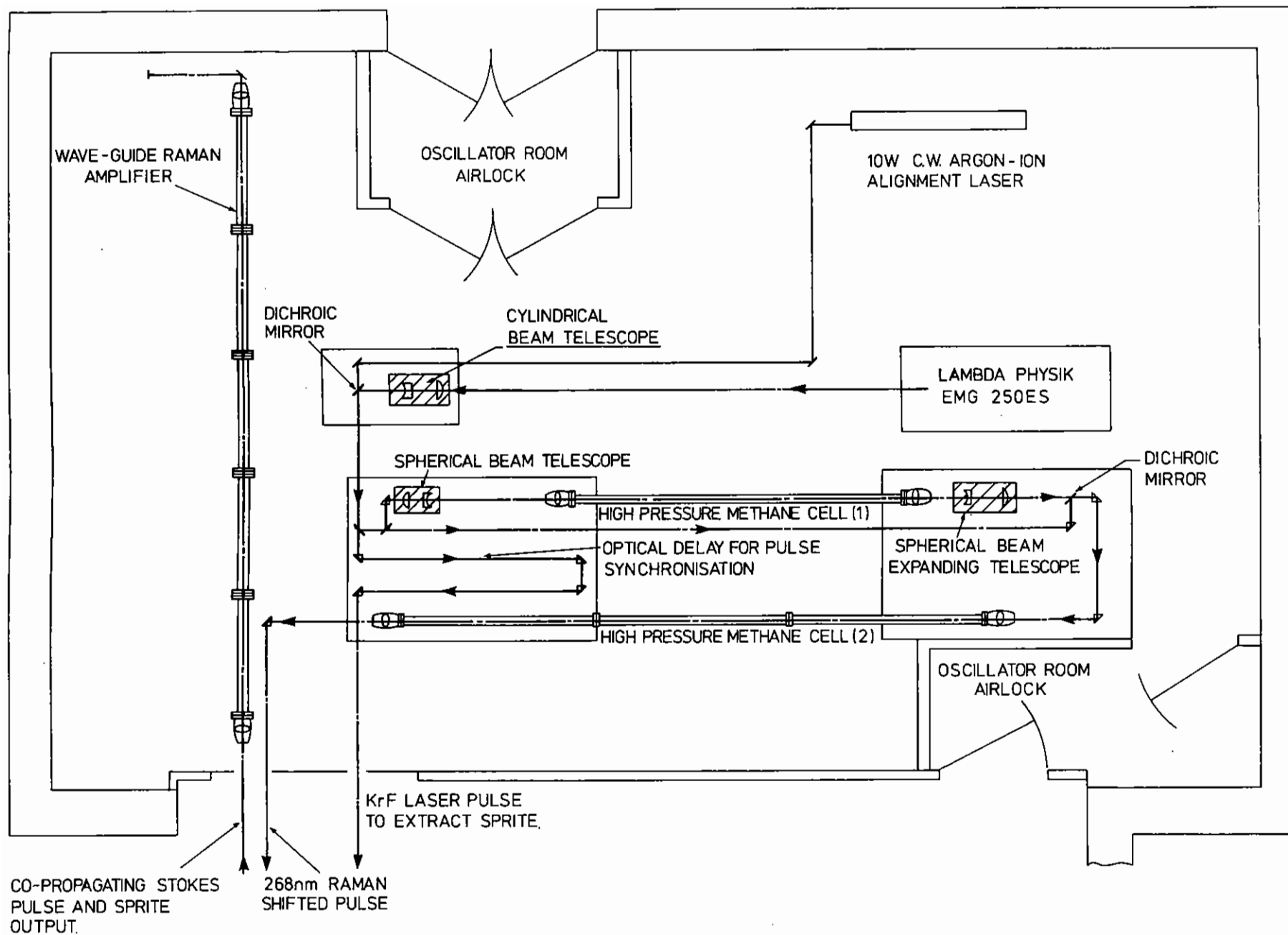


Fig. 2.11 Raman oscillator room.

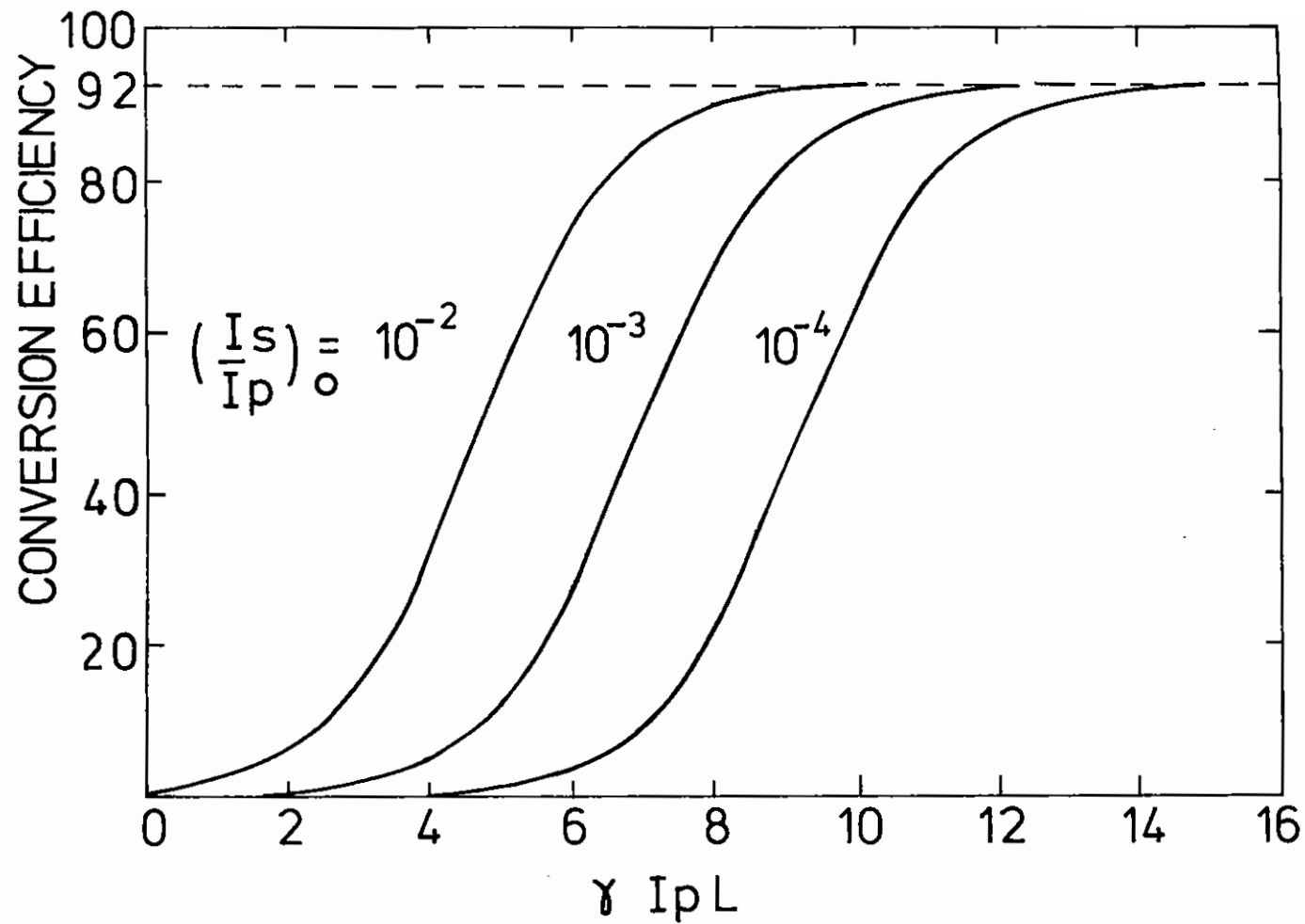


Fig. 2.13 Theoretical conversion efficiency of a forward Raman amplifier as a function of Stokes input intensity. Calculation assumes total beam overlap and no Stokes or Pump loss processes.

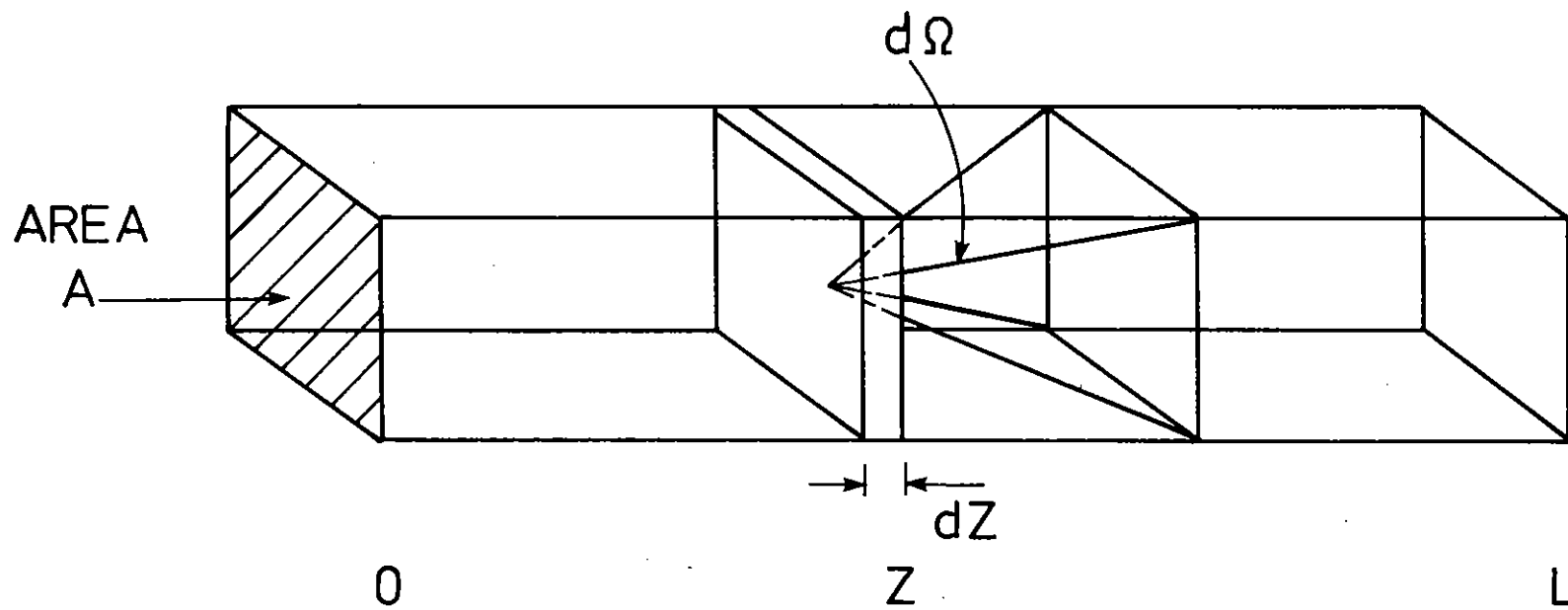


Fig. 2.14 Geometry for calculating amplified scattered Stokes intensity in a light guide.

of the multiple plane wave pumping and a rectangular guide geometry is that this intensity remains constant in the light guide (apart from a small attenuation due to reflection loss from the light guide walls). The pump is nowhere brought to a focus and thus the possibility of parasitic onset is greatly reduced. Because of the large number of pump beams employed and their intensity averaging in the guide, the transverse intensity variation of the pump $I_p(x,y)$ should be very small. This implies preservation of the beam profile of the Stokes wave, $I_s(x,y)$ since:

$$\frac{\partial I_s(x,y)}{\partial z} = \gamma I_p(x,y) I_s(x,y) \quad (1)$$

where γ is the gain coefficient per unit length, per unit pump intensity.

The Raman gain γ , is given by:

$$\gamma = \frac{\lambda^2}{h\nu} \frac{N_s}{\pi \Delta \nu_r} \frac{\partial \sigma}{\partial \Omega} \quad (2)$$

where N_s is the scatterer density, $\frac{\partial \sigma}{\partial \Omega}$ the differential scattering cross section and $\Delta \nu_r$ is the HWHM of the Raman line. In methane at pressures of over one atmosphere, the linewidth is that of the rotationally broadened Q branch and thus γ depends on pressure only through N . (Unlike H_2 where pressure broadening - or narrowing - is important). Putting in values the gain coefficient in CH_4 is:

$$\gamma = 0.1 \text{ cm}^{-1} (\text{GW cm}^{-2})^{-1} (\text{atmos})^{-1}$$

The light guide length in the experiment is 2 m and the maximum pressure is 3 atm thus a maximum value of $\gamma I_p L$ of 30 may be obtained.

The required value of $\gamma I_p L$ is determined by the input intensity and the need to deplete the pump. The solution of (1) (and its coupled equation for pump photons) under conditions of pump depletion is:

$$I_s(L) = I_s(0) \left(\frac{\omega_s}{\omega_p} \right) \frac{(I_s/I_p)_0 e^{\gamma I_p(0)L}}{(I_s/I_p)_0 e^{\gamma I_p(0)L} + 1}$$

where $(I_s/I_p)_0$ is the input intensity of the Stokes wave relative to the pump intensity. This expression is plotted in Figure 2.13 for the case of KrF pumped methane amplifiers when the maximum conversion efficiency is limited by the quantum ratio to 92%. Note this calculation considers no other pump or Stokes loss mechanisms and assumes complete beam overlap, however it does show that for a single-pass amplifier it is possible in principle to simultaneously achieve both high gain ($> 10^3$) and high conversion efficiency, a property which is not possessed by conventional laser amplifiers.

The useful upper limit of gain will be determined by pump depletion due to amplifier non-stimulated Raman scattering, a process akin to ASE in normal amplifiers. An estimate of this limit may be obtained as follows. From the definition of the scattering cross section, $\frac{d\sigma}{d\Omega}$ the power scattered into a solid angle $d\Omega$ from a volume element dV is:

$$dP_s = I_p N_s dV \left(\frac{d\sigma}{d\Omega} \right) d\Omega$$

With reference to Figure 2.14 the amplified scattered intensity at L due to the scattering from the volume element at Z is

$$AdI_s = I_p N_s Adz \frac{d\sigma}{d\Omega} d\Omega e^{\gamma I_p(L-Z)}$$

which may be simply integrated between 0 and L since $d\Omega$ is a constant which for perfectly reflecting walls is 2π . Also we assume no pump depletion. Hence:

$$I_s(L) = N_s d\Omega \left(\frac{d\sigma}{d\Omega} \right) \frac{1}{\gamma} e^{\gamma I_p L}$$

substituting γ from (2) we get:

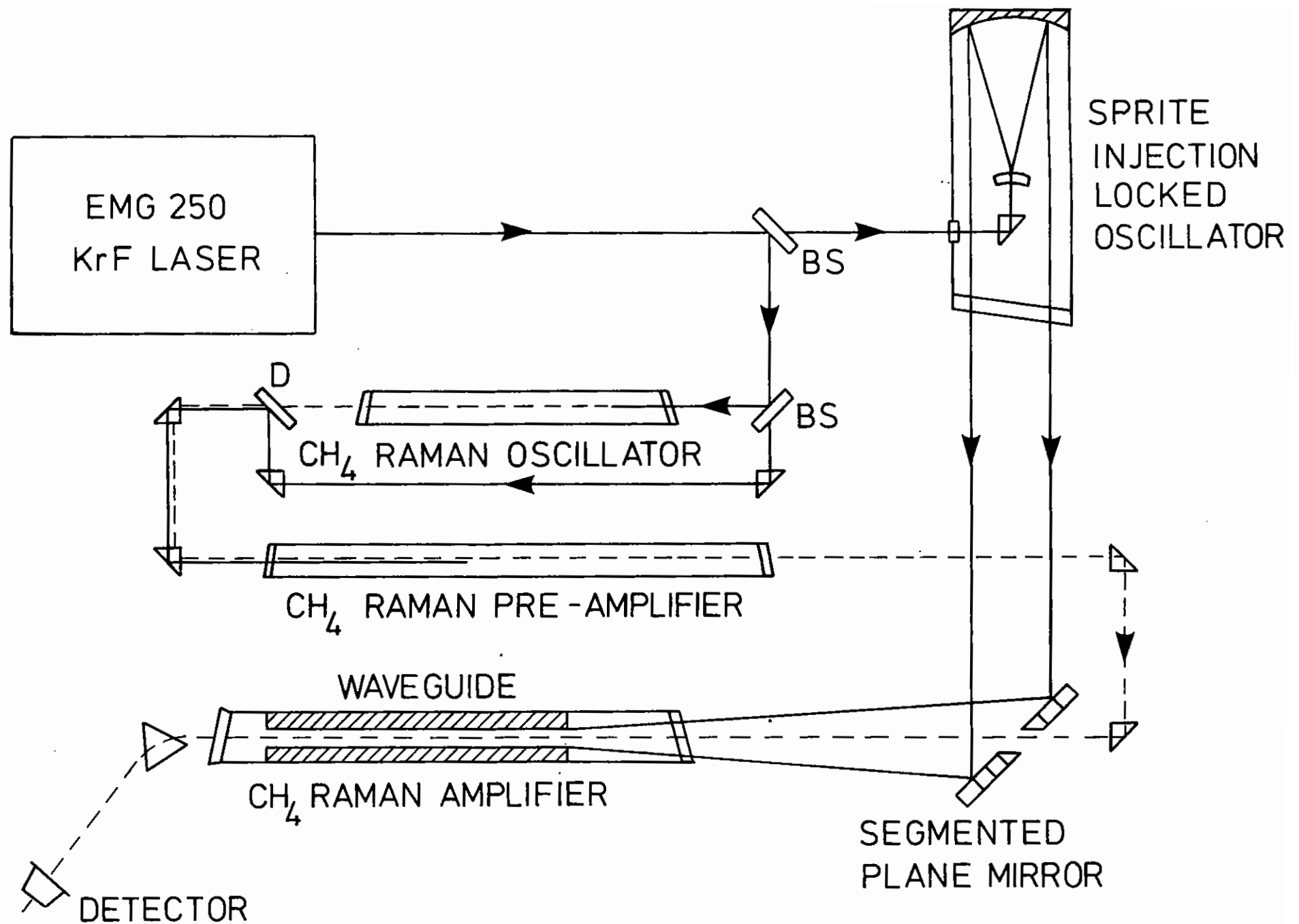


Fig. 2.12 Layout of the experiment to measure the efficiency of a high power Raman amplifier using multiple light guided pump beams. Solid lines show pump beams (248nm); dotted lines show CH₄ Stokes beams (268nm). BS = beam splitter, D - dichroic mirror.

$$I_s(L) = I_c d\Omega e^{\gamma I_p L} \quad (3)$$

where I_c is the characteristic intensity

$$I_c = \frac{h\nu}{\lambda^2} \pi \Delta\nu_r$$

and behaves as though it were the intensity, per unit solid angle of the non-stimulated scattering at the input of the amplifier. It is interesting that this intensity is independent of pump intensity and pressure and is an intrinsic property of the gain medium. Essentially it gives an input intensity below which it is impossible to obtain amplification by SRS. For methane its value is $6 \text{ W cm}^{-2} \text{ sterad}^{-1}$. Equation (3) may also be used to calculate the maximum useful value of $\gamma I_p L$ by requiring $I_s(L)$ to be less than, say 1% of I_p . For $I_p = 300 \text{ MW cm}^{-2}$ this gives $\gamma I_p \lambda < 12$ for $d\Omega = 2\pi$ ie perfectly reflecting walls and $\gamma I_p L < 18$ for $d\Omega = 10^{-3}$ sterad a value more appropriate to a glass light guide.

For the proposed experiment, the light guide is constructed from uncoated float glass. In order to obtain high transmission the pump beams are incident on the guide at angles close to grazing. The reflectivity is given by the well known Fresnel equations which for grazing angles of α rad give

$$\begin{aligned} R_s &= 1 - 2.3 \alpha \\ R_p &= 1 - 7.5 \alpha \end{aligned}$$

The maximum incident grazing angle is 20 mR which gives $R_s = 95\%$ and $R_p = 85\%$. The guide transmission is given by $T = R^n$ where n is the number of bounces. For a guide of length L , and square section of dimension a , the number of bounces is $L a/a$. For a beam element incident at grazing angle α_s to the face giving s reflection and α_p to that giving p reflection the light guide transmission is:

$$T_{\alpha_s \alpha_p} = R_{\alpha_s}^{L \alpha_s / a} R_{\alpha_p}^{L \alpha_p / a}$$

The total transmission of the guide is then calculated by averaging the transmission of the beam element over the entire range of input angles. For the proposed experiment the result of such calculations gives $T = 85\%$. Polarization effects other than the s and p reflectivities above have not been included in this analysis since for these small grazing angles the reduction in conversion efficiency due to out of plane polarization is at most a few per cent.

The experiment shown in Figure 2.12 is now being set up and will be reported on in next year's report.

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2.3 KrF Laser Development

2.3.1 Large Aperture, Long Pulse KrF Discharge Laser Feasibility Study

2.3.1(a) Introduction

The work reported here forms part of a design study for an X-ray preionized discharge-pumped RGH laser amplifier. The proposed application of the device was to serve as the final stage in the chain of preamplifiers for the Sprite pulse compression project, in which it was required to deliver 15 J in a flat-topped pulse of 100 ns when operated on the KrF transition at 249 nm. At this performance level it would exceed (by a factor of three) the highest energy reported for any self-sustained KrF discharge laser reported to date while maintaining a pulse duration comparable to the best that has been achieved. A critical factor in determining the overall performance of the laser is the behaviour of the electrical circuitry which controls the characteristics of the discharge. To investigate several questions concerning the design of the pulse-forming network (PFN) we have carried out a number of tests on small scale versions of the proposed device.

The PFN performs two more or less independent functions. The first is to ensure a uniform breakdown of the laser gas discharge. The second is to maintain a constant input of power to the glow discharge thus formed. Successful initiation of the discharge requires the very rapid application of a high voltage pulse - ideally with about five times the voltage needed to sustain the discharge (2.12).

If the overvoltage initially applied is too low, the breakdown time is inconveniently long and the discharge tends to become unstable. Once a uniform discharge is established the PFN must be capable of supplying the necessary power for the required duration of the pulse, and it should then switch off cleanly with as little post-pulse ringing as possible.

The conventional approach to meeting these requirements is to charge a section of transmission line to the overvoltage needed for good breakdown prior to connecting it across the discharge gap via a fast switch (2.13), so that for a given charging voltage the line impedance controls the power input to the discharge. Since the voltage at breakdown is some five times the sustaining voltage, the line impedance needs to be four times that of the sustained discharge, and the mismatch inevitably implies an inefficient use of initially stored energy. Ideally the line impedance should be matched to that of the discharge when running under steady conditions, but such a choice of line impedance gives rise to an initial open circuit voltage which is only twice the running voltage.

We have investigated methods by which a high voltage spike may be superimposed on the open circuit voltage of the line thus allowing a line properly matched for the fully developed discharge to produce rapid and uniform breakdown. We have also designed and tested a compact folded transmission line capable of driving the required high power discharge.

2.3.1(b) Low Voltage Simulations

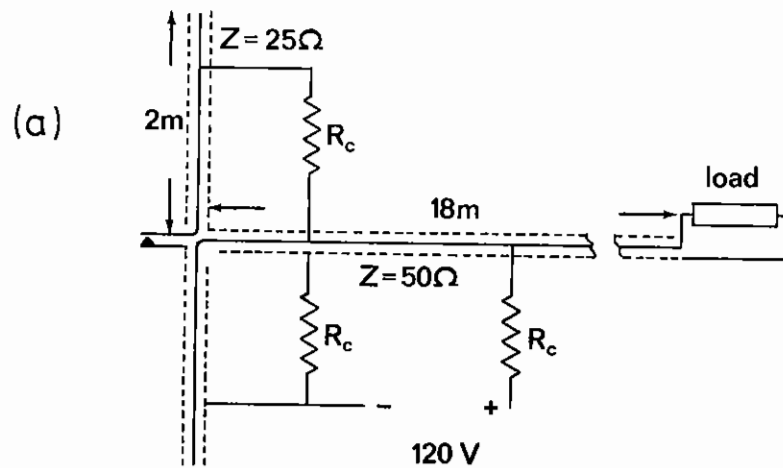
Preliminary investigations were carried out at low voltage on two types of spike-generating circuits - the "double stub" and "asymmetric Blumlein"

circuits shown in Figure 2.15. Each circuit was charged to 120 V and, via a low inductance fast acting mechanical switch, fired into a dummy load comprising a pair of 2N2222A transistors connected in series with a resistor of 18 Ω so as to provide a hold-off voltage of 210 V and a load impedance of 50 Ω when conducting.

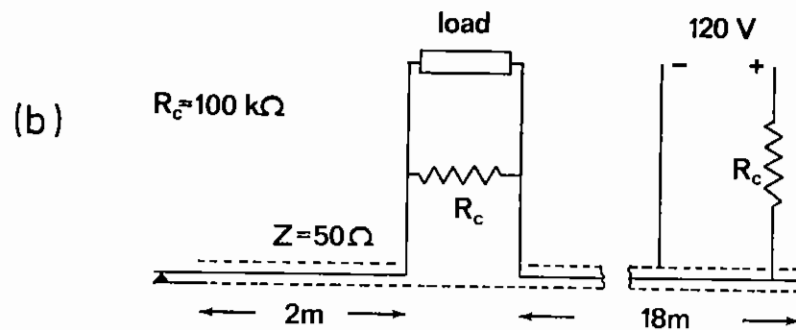
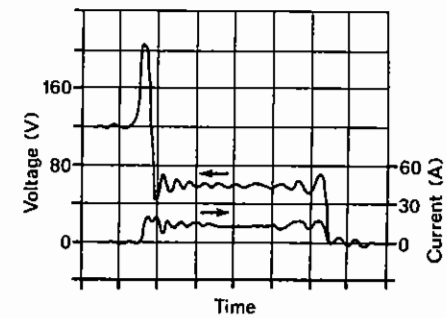
The performance of the "double stub" system was very nearly ideal (Figure 2.16(a)). The two 25 Ω stubs at one end of the line are charged in parallel and connected in series by the closing of the switch. A pulse of double the charging voltage and 20 ns duration is launched along the line and causes breakdown in the dummy laser discharge on arrival at the far end. The line then continues to discharge into the matched load giving a square current pulse with no post-pulse ringing.

In the "asymmetric Blumlein" system, closing the switch causes the voltage across the short arm of the network to invert and thus applies across the load a voltage equal to twice the charging voltage causing the dummy discharge to break down. However it is not possible to match this system to the load because the effective impedance of the line drops from 100 Ω to 50 Ω only after the initial voltage spike. Unfortunately the load impedance falls from its very high value prior to breakdown to a value near 50 Ω at an early time during the breakdown phase. Severe voltage and current ringing are therefore expected and are indeed observed in this device - a major disadvantage of an otherwise attractive circuit configuration.

At first sight, therefore, it might seem that the better of the two systems is the double-stub configuration. However, as a practical laser driver it does present some problems. Firstly, the construction of two half-impedance stubs capable of operating at the full working voltage and still sufficiently small to connect smoothly to the remainder of the line poses some practical difficulties (similar to those encountered in the folded PFN - see below). Secondly the whole PFN has to withstand an (albeit short) pulse of double the charging voltage propagating along it. Finally, the appearance of the charging voltage across the load during the pulse charging cycle would probably require that the laser should be



(a)



(b)

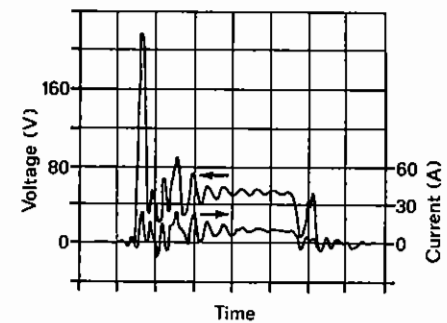


Fig. 2.15 (a) "Double - stub" Pulse Forming Network
(b) "Asymmetric Blumlein" Pulse Forming Network.

Fig. 2.16 Voltage and current at load for circuits shown in Fig. 2.15.
(a) Double - stub PFN.
(b) Asymmetric Blumlein PFN
Time - scale 40ns per division.

isolated from the PFN by a fast switch. This would inevitably add to the risetime and reduce the efficiency of the system overall.

By comparison, the electrical stresses likely to be encountered with the Blumlein circuit are smaller, and recent work reported by Sze (2.14) suggests that the current modulation would have very little effect on laser output - gas kinetic mechanisms fortunately act so as to reduce variations in gain caused by ringing in the input power waveform. For these reasons it was decided to concentrate all further development effort on the asymmetric Blumlein device.

2.3.1(c) High Voltage Tests

(1) Load Devices

For the tests of PFN configurations operating at realistic values of voltage two types of load were used. The first was a "dummy laser" comprising a liquid electrolyte resistor in series with a spark gap to simulate the breakdown characteristics. The second was a specially constructed test laser chamber containing a pair of electrodes fabricated from 25 mm brass rod profiled so as to provide a transverse discharge of 6 mm width and 210 mm length with an electrode spacing of 19 mm. Ultra-violet preionization of the discharge channel was provided by a spark array on one side of the main electrodes.

(ii) Discrete Element Asymmetric Blumlein PFN

Initial high voltage tests were carried out using an asymmetric Blumlein lumped component circuit whose long section was made from 8 pairs of 2.3 nF ceramic capacitors clamped at 50 mm intervals between two aluminium plates of 120 mm width and 40 mm separation so as to provide a line of impedance 2.1 Ω and capacitance 37 nF. This line was charged in 500 ns from a DC charged storage capacitor (70 nF) via an inductor (1 μ H), as shown in Figure 2.17. The short section of the line consisted of a single pair of ceramic capacitors and a corona-preionized rail gap. Although the long section of the line approximates well the behaviour of the continuous

line, the discrete nature of the components in the short section indicates that it must be treated as a simple L-C inverter.

The spark gap switch was originally designed as a simple rail gap with 12.5 mm diameter brass electrodes contoured so as to provide a gap of 6 mm spacing over a length of 60 mm and mounted in a perspex tube of 33 mm ID pressurized with air at up to 35 psig. In operation it proved to be disappointingly unreliable, displaying a large variation in firing voltage level, and a large percentage of misfires. Even when the gap was reduced to 3 mm and firing was reliable an unacceptable jitter still remained. To overcome this problem, provision was made for corona preionization of the gap by adding two overlapping sheets of aluminium foil, one attached to the bottom electrode and wrapped around the outside of the spark gap body while that attached to the top electrode ran in contact with the inside wall of spark gap body. Where the two foils overlap in the mid-plane, notches are provided in the edge of the inside foil (Figure 2.18). When voltage is applied to the top electrode the insulating inside surface of the spark gap body charges by drawing corona from the notched foil edge, and in doing so, irradiates the gap with deep UV, and so providing preionisation.

The provision of this simple addition dramatically improved the performance of the gap. Although the voltage hold off was reduced somewhat (necessitating a return to the original 6 mm spacing) the shot-to-shot reproducibility was excellent and the spark gap conducted reliably in five channels, one opposite each of the notches cut in the inner foil. The rise time of the voltage spike produced when the gap closed was 20 - 25 ns - comparable to the pulse length per circuit element of the long section of the line.

(iii) Discrete Element PFN - Laser Tests

Typical voltage and current traces recorded with the discrete element asymmetric Blumlein PFN feeding the test laser as shown in Figure 2.19(a). Several features of the traces were strongly dependent on gas composition (He:Kr:F₂) but the voltage rise time of 30 - 35 ns (limited by the time

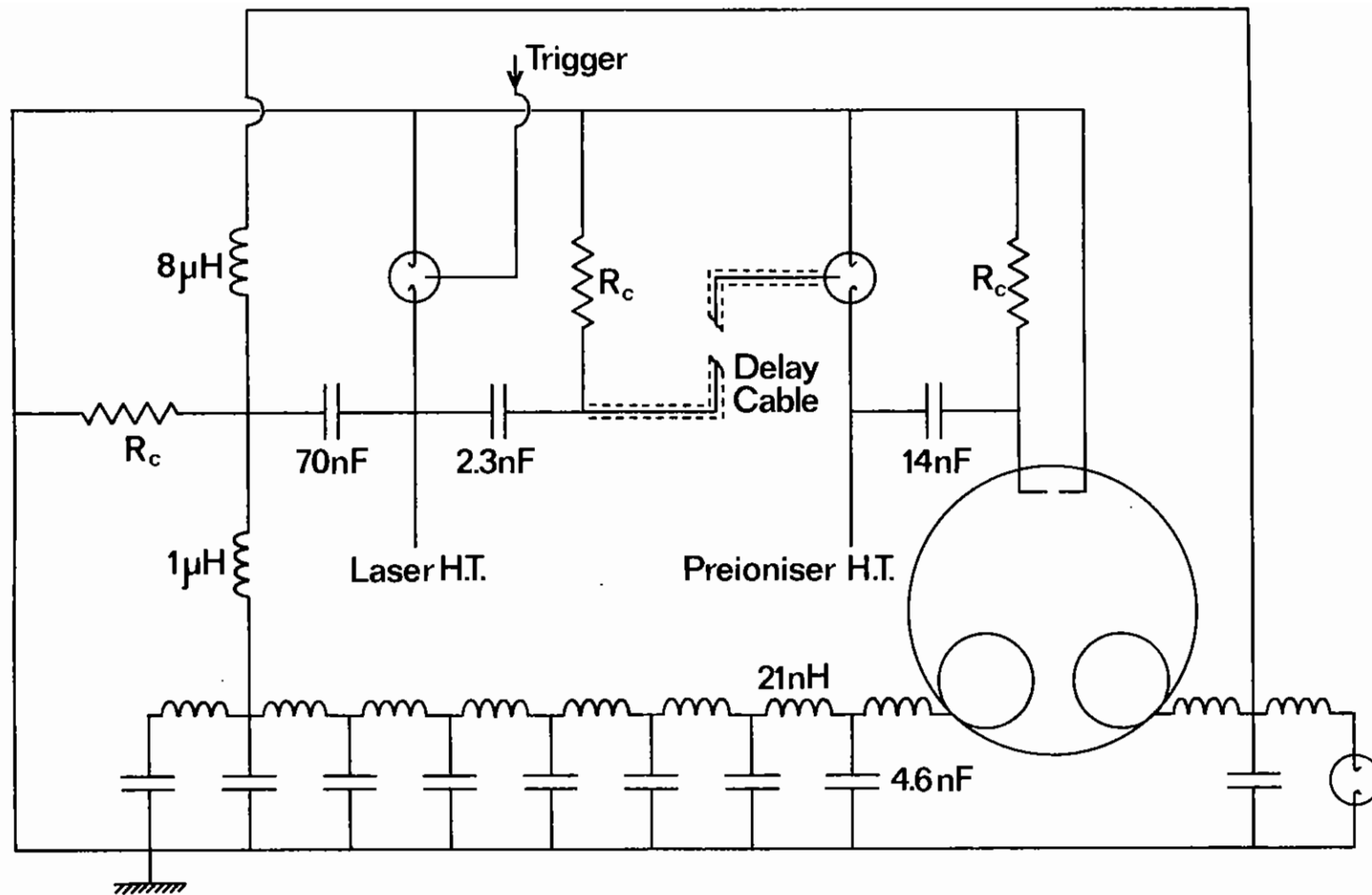


Fig. 2.17 Lumped - component asymmetric Blumlein PFN and associated pulse - charging and preionized rail gap plasma switch.

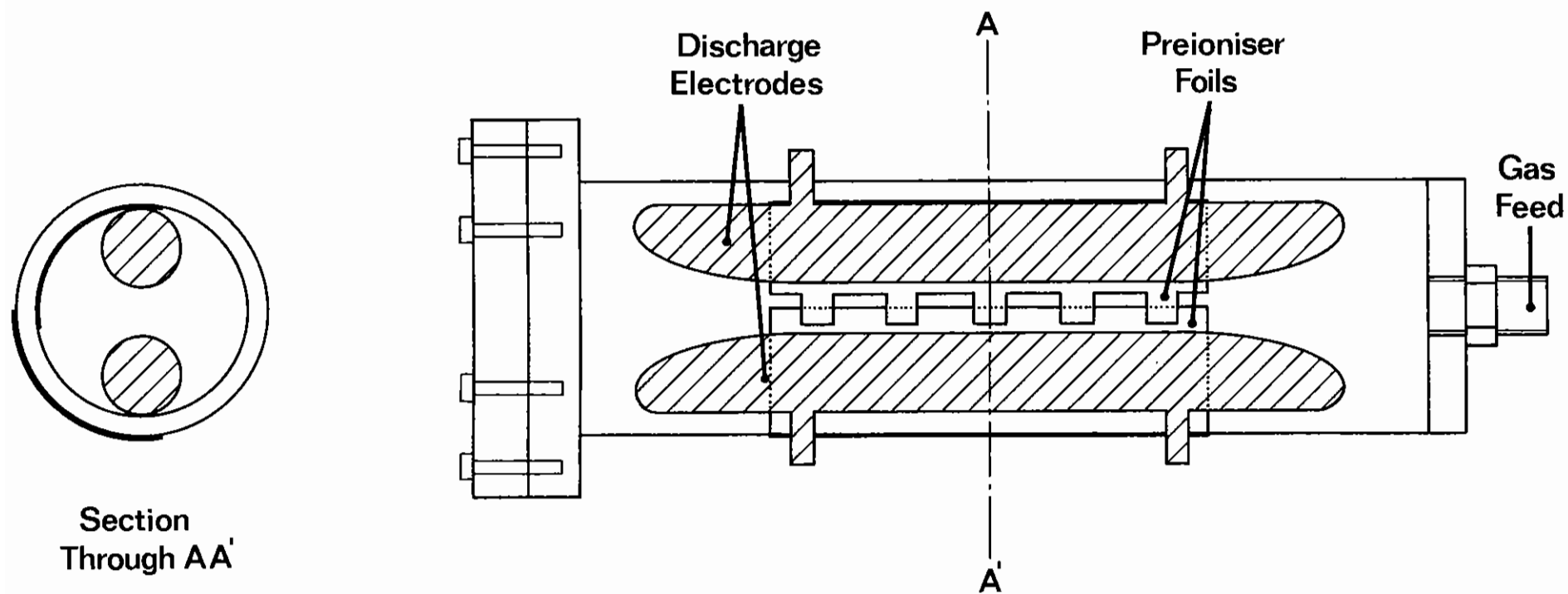


Fig. 2.18 Corona - preionized rail gap plasma switch.

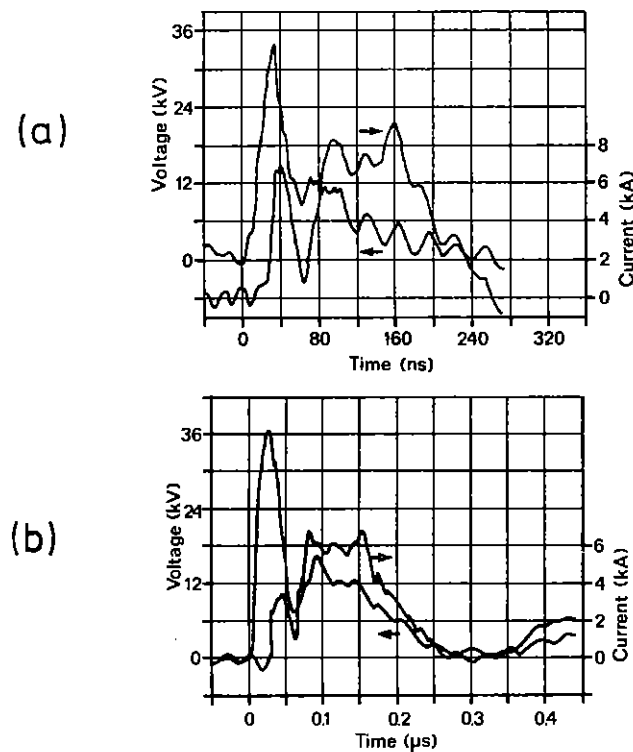


Fig. 2.19 Current and voltage at load for the circuit shown in Fig. 2.17.
 (a) Coupled to KrF discharge laser.
 (b) Coupled to dummy laser load.
 Note change of time scale between (a) and (b).

taken for the short section to invert voltage via the rail gap switch) is a feature common to all. Because this rate of voltage rise is slower than ideal the laser gap breaks down at an overvoltage of only 3 times the expected running voltage. After taking 20 - 30 ns to recover from the initial breakdown transient the current is more or less stable for the following 80 ns. About 70 ns after breakdown the voltage across the laser gap collapses - an event which occurs at progressively earlier times the greater the proportion of F_2 in the laser mixture. Laser action at 249 nm began during the initial current spike. Although 249 nm spontaneous emission persisted a little longer, termination of laser action appears to coincide with voltage collapse and is almost certainly associated with the onset of the same kind of constriction instability which has been observed in studies of the pulse length limitations of XeCl lasers (2.15). The longest laser pulse duration observed was 50 ns (base width).

In order to establish which features were related to the PFN design (and not artifacts due to imperfectly profiled electrodes or non-uniform preionisation density of the test laser) the PFN was tested with the dummy laser load. The resulting traces are shown in Figure 2.19(b). It is noticeable that the early voltage collapse is no longer evident and that the voltage is maintained at a sensibly high value for about 125 ns, and it may be safely concluded that the PFN was operating in the desired manner.

2.3.1(d) Design and Testing of a Compact Transmission Line PFN

(i) Design Considerations

Storage of energies in the kJ range needed for the full scale device in PFN's employing solid dielectric media presents severe practical problems and a much more attractive approach is to take advantage of the high dielectric constant of water in the design of the PFN. Since the PFN is required to deliver pulses of 48 kV and 150 ns duration into a load of 0.5Ω and to have an open circuit voltage higher than 200 kV when operating as a Blumlein network, it should have an impedance of 0.5Ω and be capable of working at charging voltages of about 100 kV.

For an electrical pulse length of 150 ns the physical length of the long section of the asymmetric Blumlein is 2.5 m. In order to keep the overall dimensions of the PFN unit as compact as possible this element appears in folded form in the final design shown in Figure 2.20. The line comprises two sets of interdigitating fins separated by the water dielectric and perspex spacers. The spacing d , between opposing fins is related to w , the width of the line by

$$d = \sqrt{\frac{\epsilon\epsilon_0}{\mu\mu_0}} Z w = \frac{Zw}{42}$$

in which Z is the line impedance in ohms. For a line of given Z , the larger the value of w the larger can be the value of d , thus reducing electrical stress in the dielectric. Design constraints on the dimensions of the final laser device restrict w to 900 mm, so that for $Z = 0.5 \Omega$ the value of d should be 10.8 mm. At a charging voltage of 100 kV this implies a field strength of 93 kV cm^{-1} in the uniform field region - a value well within the breakdown strength of water (160 kV cm^{-1} for the appropriate area and pulse length (2.16)). However the maximum electrical stress is expected to occur at the bends of the line where it is folded since the enhancement factor f by which the field at the tip of a fin may exceed the uniform region value will be considerable here. Approximating the fin edge and opposing plate surface by two co-axial half cylinders of radius r and $(r+d)$ respectively, the value of f is given by

$$f = \frac{d}{r \ln(1+d/r)}$$

To allow for a safety margin of 25% the value of f must not exceed 1.3 which for a d of 10.8 mm implies $r > 17 \text{ mm}$. Thus the minimum thickness of the material from which the fins are fabricated is 34 mm - implying an inconveniently massive assembly if the fins are machined from sheet aluminium alloy. In practical terms a thickness of alloy sheet of 6.25 mm is the best choice, but the enhancement factor f that this choice of r implies is 2.4, which would restrict the charging voltage of the line to

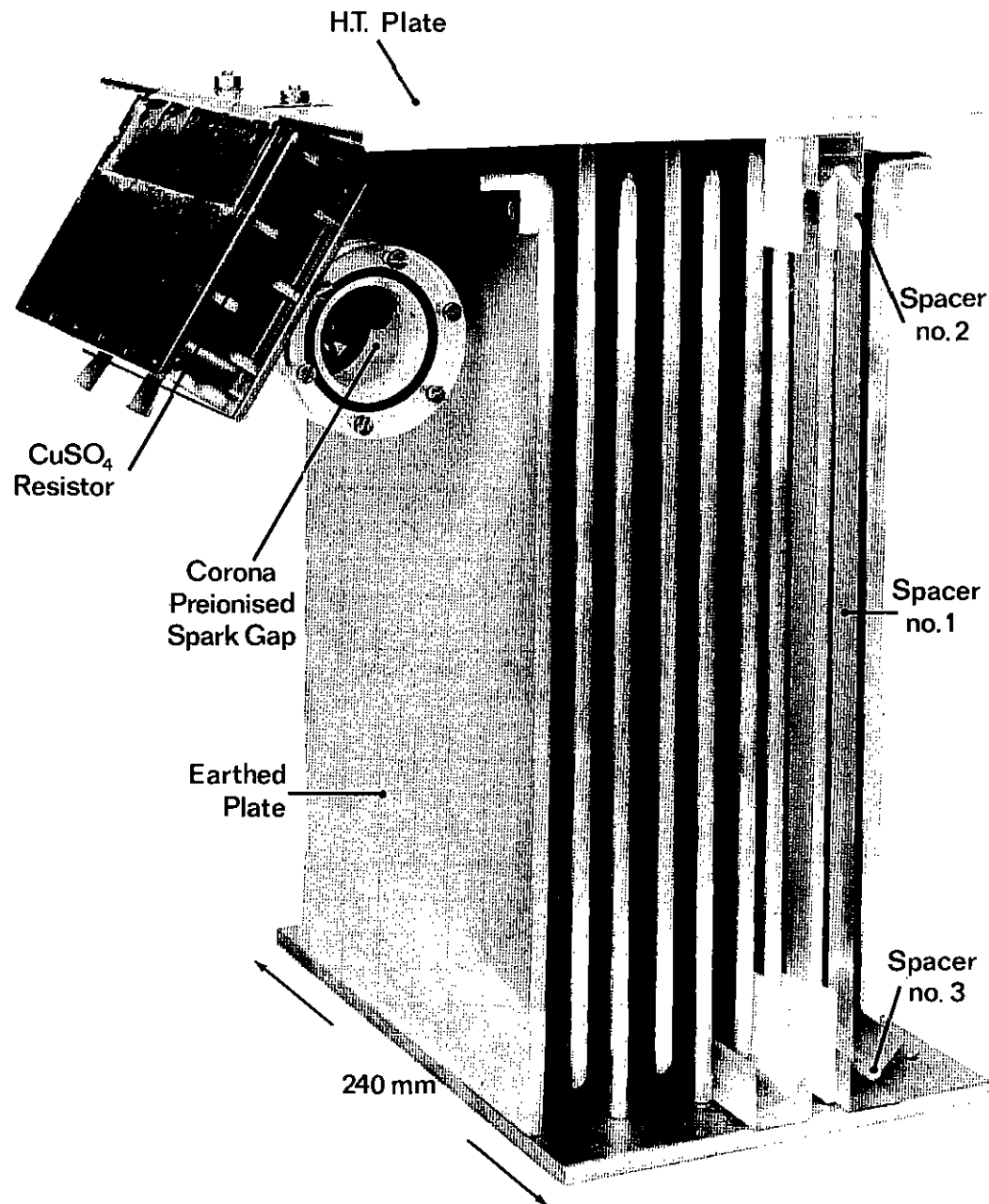


Fig. 2.20

Photograph of folded water line assembly with dummy laser load attached.

an undesirably low value of 70 kV. It was therefore imperative to find a method of protecting the line from breakdown at the fin tips while keeping the fin thickness to 6.25 mm.

(ii) Tests of Compact Water Filled Line Without Fin-edge Protection

The folded transmission line device shown in Figure 2.20 is a one quarter width model of the long section of the proposed asymmetric Blumlein PFN, with a slightly reduced spacing ($d = 9.5$ mm) but with the full electrical length. When immersed in water the total capacitance of the line is 38 nF and its impedance of 2Ω is four times as large as that of a full width version. The conductivity of the water bath is maintained at a low value by continuous recirculation of the water through an ion-exchange column.

For tests, the line was connected to a storage capacitor from which it was charged to 60 kV with no evidence of breakdown. The peak electric field in the unprotected fin edge regions was 145 kV cm^{-1} , sustained over small areas for short periods. The voltage and current pulses observed when the line was discharged into the dummy laser load were square topped for the first 120 ns but showed evidence of droop during the final 30 ns - decreasing to 70% of the plateau value by the end of the pulse. It is possible that the droop is caused by partial "short circuit" effects due to fringing fields built up at the outer edges of the fins which may cause some energy to transfer across folds in the line. If this is the correct interpretation, the effects should be much reduced on a line of full width. An alternative explanation is that the rear section of line is starting to transfer energy back into the storage capacitor by the end of the pulse. Although numerical analysis suggests that this explanation is not a likely one, it was found that the current droop could be reduced by firing the dummy laser gap before pulse charging of the PFN was complete. The consequent drop in initial voltage could be compensated by increasing the charging voltage on the storage bank.

(iii) Tests of Compact Water-filled Line with Fin-edge Protection

To overcome the problem associated with breakdown in the fin-edge regions,

a set of perspex spacers was constructed and installed (see Figure 2.20). Perspex is a material of high dielectric strength (2 MV cm^{-1}) and its relatively low dielectric constant ($\epsilon_r \approx 3$) results in a field distribution which confines the regions of maximum stress to those containing perspex rather than water (2.17). The problem of breakdown in the bulk volume is therefore removed, but one of tracking along the perspex-water interface is substituted. It was such tracking which was to prove the ultimate limitation to the working voltage of the device.

The assembly of the perspex spacers are shown in somewhat greater detail in Figure 2.21. Spacer number 1 is a vertical strip separating the outer edges of adjacent fins, with a bevelled inner edge to increase the length of the tracking path. Spacer number 3 is a vee-block with a groove machined at the apex to accept the tip of the HV fin (shown in dotted outline in Figure 2.21), thereby separating the HV fin from the earthed plate beneath. Spacer number 2 is a later addition to the initial design, included in order to prevent tracking along the triple interface between the water and spacers numbers 1 and 3.

Since only two sets of spacers were available, the length of the line was reduced from its original 8 sections to just the two containing the spacers by removing fins in the unused sections. It was charged via a two-stage Marx generator capable of delivering 94 kV from an erected capacitance of 17.5 nF. The two section line was pulse charged to 125 kV in 300 ns.

Before adding spacer number 2, tracking occurred at a line voltage of 110 kV. With spacer number 2 in place the line survived repeated shots at 115 - 120 kV despite prolonged ringing due to the absence of any dissipative load. The line finally failed by tracking along the bevelled edge of spacer number 3 on the second cycle of a ringing waveform of peak applied voltage 125 kV. There was some evidence that trapped air bubbles had contributed to this failure. In a practical version of this device it would be desirable to install water circulation channels to flush out any impurities which might otherwise increase the conductivity of the water as well as any bubbles which would lower its dielectric strength.

2.3.1(e) Conclusions

We have demonstrated that it is feasible to construct a PFN capable of providing a waveform with a high voltage spike at the leading edge followed by a long flat-topped pulse by employing the concept of an asymmetric Blumlein circuit. The use of dummy laser loads which employ a spark gap switch to simulate the breakdown threshold of the laser gap and a liquid electrolyte resistor to simulate the dissipation of the laser glow discharge provides an excellent means for testing such PFN's free from interpretational uncertainties introduced by the discharge instabilities of imperfectly preionized transverse discharges. Finally, we have demonstrated a very compact folded water filled line suitable for driving high energy discharge lasers, and explored the limits of its operating range. With this information it is now possible to undertake the design of a PFN for the full scale laser device which would have the desired characteristics.

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2.3.2 UV Laser Damage Studies

The establishment of a laser damage facility for work in the UV region corresponding to excimer wavelengths (193 - 351 nm) was briefly described in the 1982 Annual Report. The major feature of the programme has been to ensure that the spot size dependence of the damage threshold has been removed by the use of relatively large spot sizes. Irradiation areas of $\sim 2 \text{ mm}^2$ with corresponding energy densities up to 45 Jcm^{-2} are available. This figure is 6 times the highest damage threshold for a multilayer dielectric mirror that has been measured. Characterisation of the focused beam has received careful attention using a number of techniques including; burn patterns, dylux film activation and photographic recording of the laser profile in a weakly fluorescing dye. The shot to shot variation in the output energy and beam profile has been shown to be stable to better than 5%.

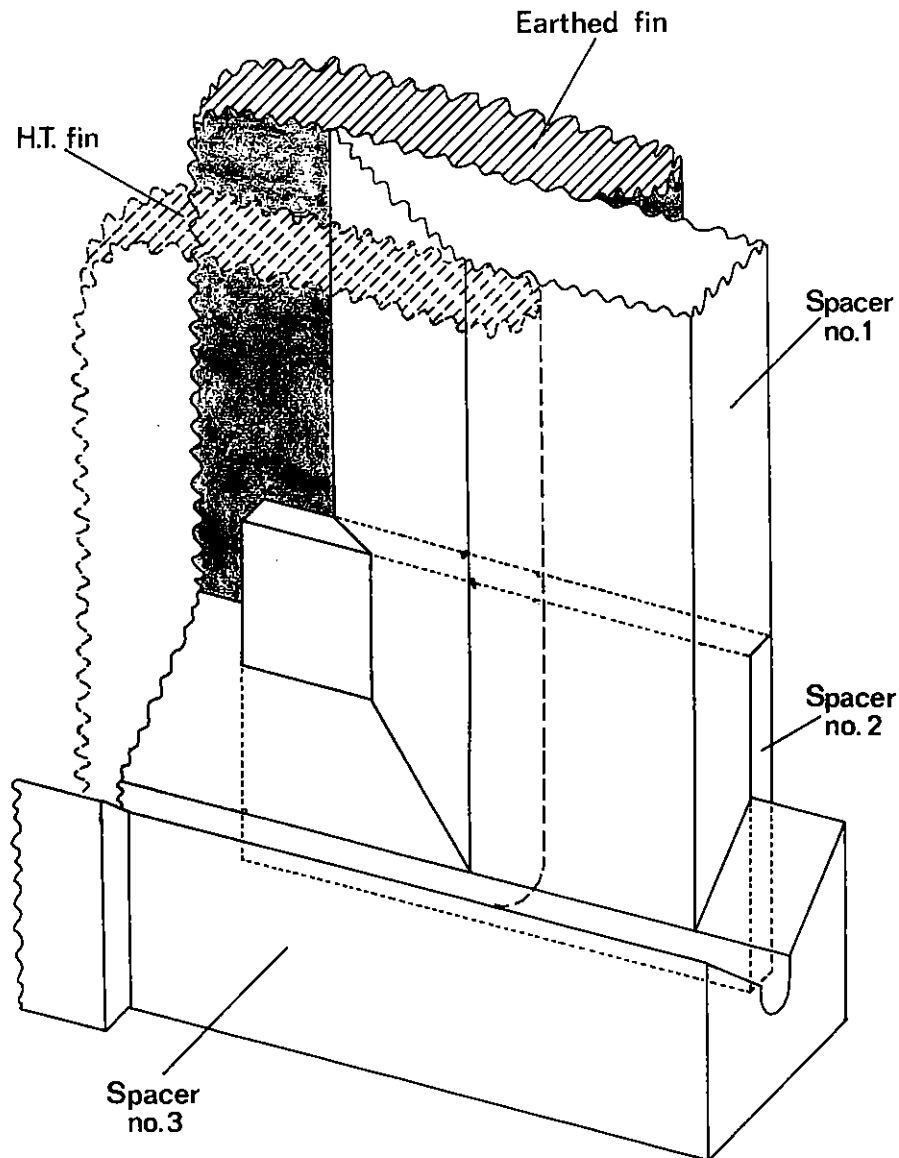


Fig. 2.21 Detail of fin-edge protecting perspex spacers installed in the folded line PFN.

There exist many different interpretations of what constitutes damage and how the damage threshold is defined. The policy adopted here with regards to handling and cleanliness is to ensure the removal of obvious dust and loose matter by use of a camera lens blower. The surface to be tested is observed with a low power (10 x) microscope, adjusted to give maximum contrast in the image plane. The damage site is observed as the laser is fired. Damage is defined as any physical change in the sample observed through the microscope. Two damage thresholds are defined corresponding to that which causes damage in every single shot and that which causes no damage in 600 shots on the same site. The single shot criterion is a measure of the maximum energy density that the film may withstand while the multishot value is perhaps a more useful practical figure. Damage studies have been made at 193 nm (ArF) on a number of alkali and alkali-earth halides. These are shown in Table 2.01. The single shot threshold values varied in the range 1 - 4 Jcm⁻². The freshly cleaved samples of NaF and LiF have higher damage thresholds. This was attributed to the high water solubility of these materials compared to the others. The polished sample thresholds are somewhat lower than their equivalent cleaved counter-parts. This suggests that polishing material remaining in the surface is absorbing at 193 nm and initiating the damage.

The majority of the work has been at 248 nm (KrF) as an aid to the development of mirrors and coatings for use in the RAL laser programme. A range of high reflectivity (nominally 99%) multilayer dielectric film mirrors has been tested at 248 nm. The results are summarised in Table 2.02. Damage tests were performed on a number of used mirrors, these being ones which have previously operated in a fluorine environment as excimer mirrors. The multishot threshold of these mirrors was found to be consistently 3 - 4 times less than the single shot threshold. The damage threshold of mirror B was previously determined by the Lawrence Livermore National Laboratory. The damage threshold of 7.9 Jcm⁻² compares favourably with the 6 Jcm⁻² measured by LLNL. Four main points arise from Table 2.02.

CRYSTAL	TYPE	SINGLE SHOT DAMAGE THRESHOLD/Jcm ⁻² + 30%			MULTI-SHOT DAMAGE THRESHOLD/Jcm ⁻² + 30%		
		No 1	No 2	No 3	No 1	No 2	No 3
LiF	IR Grade	2.9*	2.2	2.6			
	UV Grade	2.4	3.3				
	Polished	1.8			1.0		
NaF	Cleaved	1.3	2.9*	3.0*			
CaF ₂	IR Grade	2.1	1.9				0.5
	UV Grade	2.5	2.6			0.65	
MgF ₂	Polished	2.6			1.8		
SrF ₂	Cleaved	1.6	2.2				
	Polished	1.3			0.5		
BaF ₂	Cleaved	1.2		1.0		0.6	

TABLE 2.01 DAMAGE THRESHOLDS AT 193nm (ArF) OF UNCOATED
FLUORIDE CRYSTALS

TYPE	DESCRIPTION	SINGLE SHOT DAMAGE THRESHOLD/Jcm ⁻²	MULTISHOT DAMAGE THRESHOLD/Jcm ⁻²
A	USED	0.6	
B	UNUSED	7.9	2.4
C	UNUSED	2.6	
	UNUSED	2.0	0.5
	USED	0.3	
	USED	0.3	
	USED (DURAN SUBSTRATE)	0.75	

	FUSED SILICA SUBSTRATE (DYNASIL 4000)	12.3	4.6

TABLE 2.02 MULTI-LAYER DIELECTRICS REFLECTORS AT 249nm

- (i) Mirror B was designed for a maximum damage threshold in a non fluorine atmosphere whereas the other unused mirrors (c) all involved fluoride coatings to allow operation in this gas.
- (ii) The six-fold reduction of the damage threshold for thin film coatings compared to the uncoated substrate, confirm the damage limitations of coated optics.
- (iii) The mirrors exposed to a fluorine environment have a lower damage threshold than unused mirrors of the same kind.
- (iv) There seems to be no great advantage in using substrates of fused silica compared to BK7 or Pyrex when the mirror is a high reflector.

Large diameter multilayer dielectric coated components are difficult and expensive to produce and there may be considerable advantages to be obtained from the use of metal mirrors in the ultra-violet. Aluminium is the only metal with a high reflectivity (~ 90%) in the region of interest and a series of experiments is currently being undertaken to try to find the optimum deposition parameters for overcoated aluminium mirrors. The single shot damage thresholds for MgF₂ and ThF₄ overcoated aluminium mirrors are shown in Table 2.03 and Table 2.04 respectively.

The main conclusions are that the dielectric overcoat damages before the aluminium coating and the substrate material has little influence on the damage threshold because the damage appears to be initiated in the thin film. The aluminium mirrors presented in Table 2.04 had varying aluminium thickness and dielectric overcoating thickness to see if this influenced the damage threshold. The damage threshold increases with reflectivity. A comparison between mirrors in Table 2.04 cannot really be made because of the widely varying reflectivities. This illustrates the precise control required on the evaporation conditions to achieve maximum reflectivity in the ultra-violet.

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NUMBER	SUBSTRATE	DAMAGE THRESHOLD /Jcm ⁻²	REFLECTIVITY % @ 30° INCIDENCE
1	BK-7	.18	
2*	FUSED SILICA	.17	
3	COPPER	.2	69 ± 4
4	COPPER/NICKEL	.17	84 ± 3
<u>BACK COATED MIRRORS TESTED THROUGH QUARTZ</u>			
5	FUSED SILICA	.6	
2*	FUSED SILICA	.23	78 ± 6

* Tested in both directions

Al thickness = 1000Å

No 1 and 5 5500Å MgF₂

No 2, 3, 4 λ/2 @ 249nm OF MgF₂

TABLE 2.03 ALUMINIUM MIRRORS OVERCOATED WITH MgF₂ DAMAGE THRESHOLD
@ 249nm

NUMBER	SUBSTRATE	REFLECTIVITY/%	DAMAGE THRESHOLD /Jcm ⁻²
1	FUSED SILICA	85	.38
2	PYREX	85	.4
3	PYREX	47	.2
4	BK-7	58	< .15
<u>TESTED THROUGH QUARTZ</u>			
5	FUSED SILICA	64	.3
6	FUSED SILICA	5 (through film) 40 (through quartz)	< .08

COATING DETAILS

1 and 2 2μm Al + λ/2 @ 249nm ThF₄

3 and 5 5000Å Al + 3λ @ 249nm ThF₄

4 1000Å Al + 300Å ThF₄

TABLE 2.04 ALUMINIUM MIRRORS OVERCOATED WITH ThF₄

2.3.3 UV Phase Conjugation

2.3.3(a) Experiments with a KrF Laser Amplifier using Phase Conjugate Brillouin Mirrors

Several experiments have been carried out which demonstrate the utility and versatility of the KrF phase conjugate (PC) stimulated Brillouin scattering (SBS) mirror which we were the first to demonstrate last year (2.18). The experimental set up has been described previously (2.18) and consists of a KrF oscillator, which produces $\sim 4 \mu\text{J}$ of energy at 248 nm in a diffraction-limited 5 ns FWHM pulse of area 0.3 cm^2 with a narrowed linewidth of $< 0.2 \text{ cm}^{-1}$, and a synchronized 46 cm x 2 cm x 0.5 cm discharge-excited amplifier module, which exhibits a small signal gain of $\sim 0.2 \text{ cm}^{-1}$ for $\sim 20 \text{ ns}$. After the first passage through the amplifier, the radiation had a divergence of $\sim 5 \times$ the diffraction limit. Although ethanol was predominantly used as the Brillouin medium in our previous work (2.18), hexane (BDH; Spectrosol grade; absorption, $< 2\% \text{ cm}^{-1}$ at 248 nm) was found to produce $\sim 30\%$ higher PC reflectivity and so was used in the present study. Highest-quality PC reflections were obtained when care was taken to ensure that the oscillator was fully line narrowed throughout the duration of the pulse.

Using different focal-length lenses to focus the radiation emergent from the amplifier into the scattering cell, in Figure 2.22 we show measurements that characterize the PC reflectivity of the Brillouin mirror as a function of the product $I f^2$, where I is the intensity at focus. As described in (2.18) spatial filtering was used to detect the PC component of the SBS return. It can be seen that the maximum observed PC reflectivity was $\sim 80\%$ and, decreases at higher intensities for a given focal length lens. Also in Figure 2.22, as expected, one can see that the reflectivity is independent of the focal length of the lens (if l is the interaction length, then the SBS signal depends on the product $I l \propto I f^2$, which is independent of f).

Because of the inability of rare-gas-halide (RGH) molecules to store

energy for times longer than a few nanoseconds, it is necessary to use various pulse compression schemes (2.04) for efficient short-pulse extraction from high-energy RGH amplifiers. One such scheme is angular multiplexing, whereby the available energy in the amplifier is extracted by a sequence of angle-encoded pulses. Efficient energy extraction from KrF using conventional optics has already been demonstrated in a $3 \times$ multiplexing experiment (2.04). A scheme for angular multiplexing using a phase conjugate mirror in a double-pass amplifier arrangement is shown in Figure 2.23. Use of a PC SBS mirror in this configuration could have several advantages over use of a conventional mirror. As well as being simple and not requiring high-quality optical components, the PC mirror has also been shown to compensate for aberrations induced on the wave front by the KrF amplifier itself (2.18). In addition, since PC mirrors automatically retroreflect beams incident upon them, they should also help to reduce cross talk between beams. As we will show later, because this type of PC mirror acts as a narrow bandwidth reflector, it also increases optical isolation of the system from broadband KrF amplified spontaneous emission (ASE). Since this type of PC SBS mirror produces a polarization state of the reflected beam similar to that obtained from a conventional mirror, ie the polarization is not conjugated (2.19), switch-out of the pulses returning from the amplifier to optical delay lines may be achieved by using a combination of polarizers and quarter-wave plates.

In Figure 2.23(a)-(e) we show oscilloscope traces of the pulses using this scheme in a $3 \times$ multiplexing arrangement (without switch-out). The diffraction limited, line narrowed oscillator pulse shown in Figure 2.23(a) is angle encoded (Figure 2.23(b)) and passed through the amplifier (Figure 2.23(c)) with a net gain of $\sim 10^3$. The reflections from the PC SBS mirror (Figure 2.23(d)) repass through the amplifier, and the diffraction-limited return from the three delay lines is shown in Figure 2.23(e). Amplitude modulation of the overall pulse shape shown in Figures 2.23(c)-(e) is not so evident as in Figure 2.23(b) because of broadening by a factor of ~ 2 of individual oscillator pulses after the initial transit through the (small-signal gain) amplifier. However, the PC SBS mirror then shortens individual pulses by $\sim 30\%$. By measuring the energy obtained from the amplifier when it is operated as a superradiant

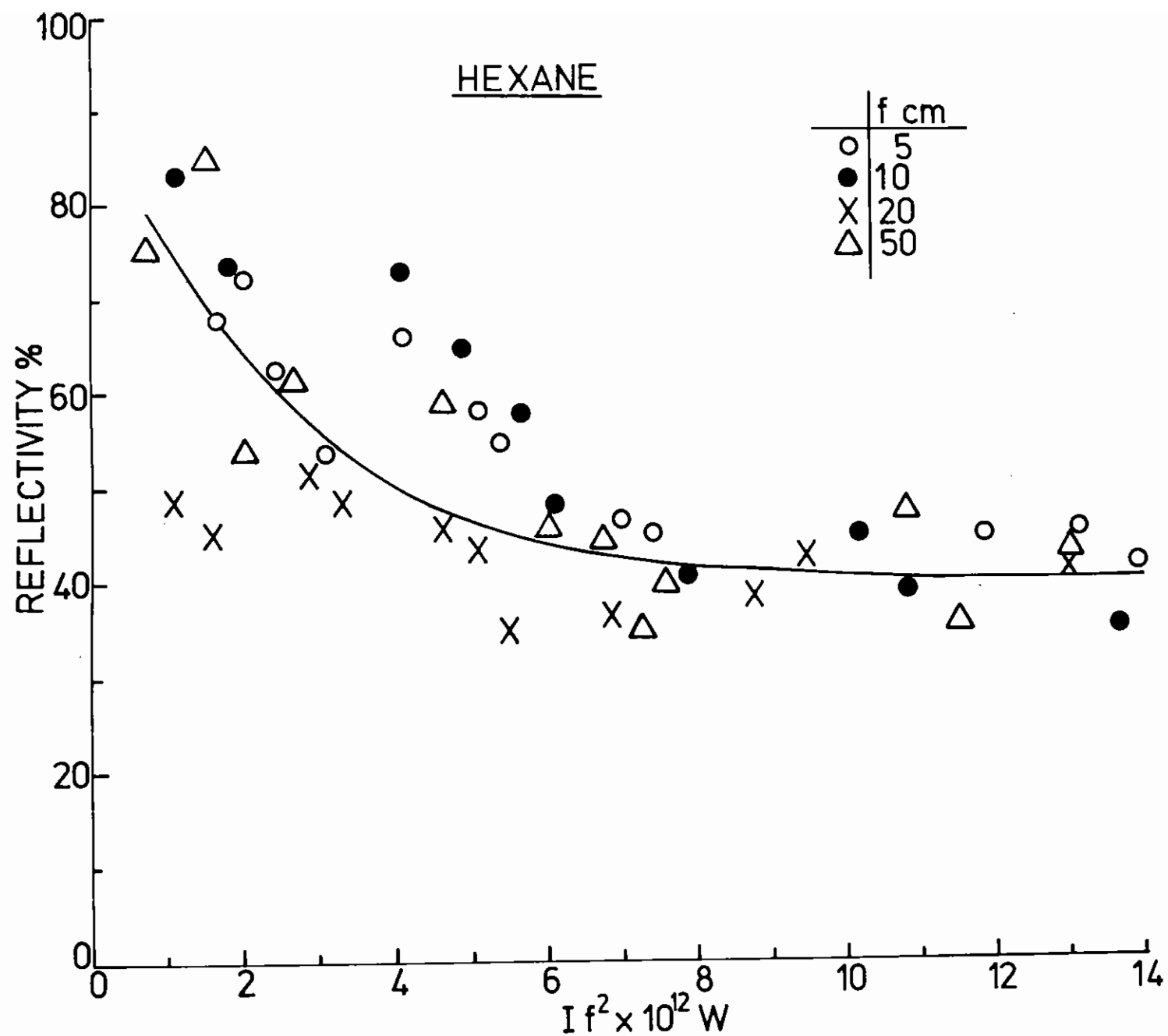


Fig. 2.22 PC SBS reflectivity versus $I f^2$ for different focal-length lenses f using a hexane scatterer.

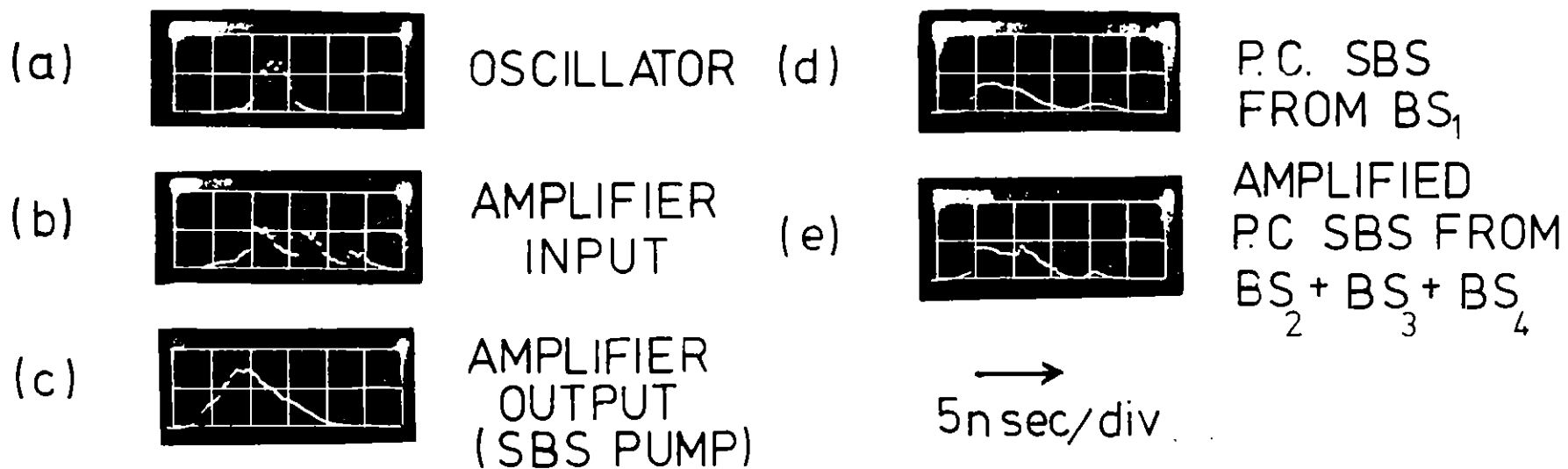
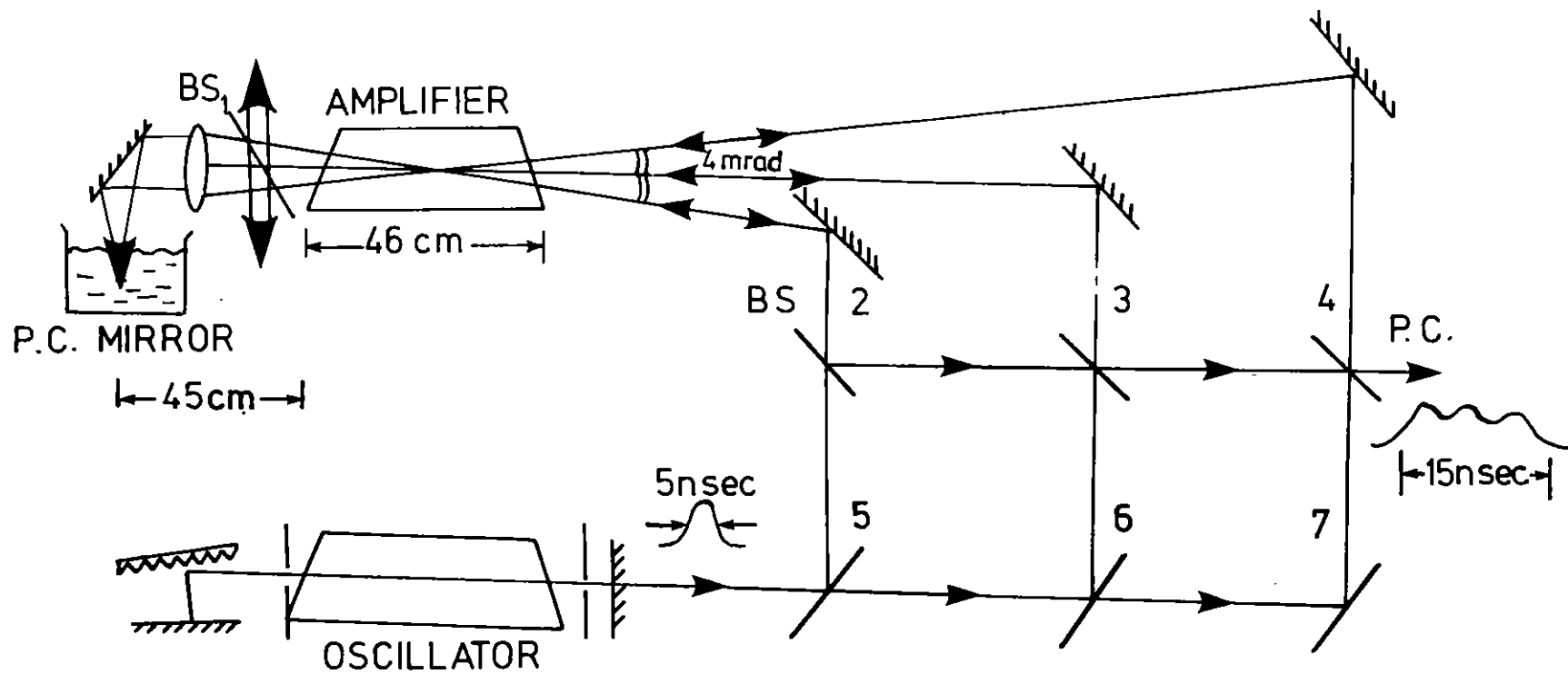


Fig. 2.23

Angular multiplexing using an SBS PC retroreflector. Each leg has 6-nsec optical delay time. BS₁-BS₄, BS₅, BS₆, and BS₇ are 10, 50, 80, and 100% reflecting beam splitters, respectively.

oscillator with an 80% plane reflecting mirror of equivalent aperture in place of the PC mirror, we were able to assess the efficiency of this multiplexing arrangement. We found that $> 70\%$ of the energy obtained from this long pulse oscillator could be extracted in the short pulse multiplexed mode of operation. In addition, the output when the SBS PC retroreflector was used was found to have a diffraction limited beam quality.

Another application of PC mirrors that we have investigated is their use as reflectors in a KrF regenerative amplifier. Regenerative amplifiers can be used to provide synchronized short diagnostic probe pulses in laser-fusion experiments (2.20). There are several advantages of using PC mirrors instead of conventional mirrors in such amplifiers. Not only are PC mirrors capable of producing an unaberrated beam, but we have found that no optical isolation of the mirrors, in the form of a Q switch, is necessary in the cavity. Also, the mirrors themselves produce pulse shortening. To demonstrate the operation of a KrF regenerative amplifier with PC mirrors, we used the set up shown in Figure 2.24. The low intensity injected oscillator pulse was passed through the Brillouin cell (M_2) into the regenerative amplifier. The output pulses obtained from the intracavity beam splitter are shown in Figure 2.24(a). It can be seen that the gain in the amplifier persists for three round trips of the injected pulse. The shortening of the first pulse to a 2 ns FWHM second pulse is a result of a combination of gain saturation and the reflection from the PC mirror, M_2 . In contrast, in Figure 2.24(b) we show the output when the PC mirror, M_2 , is replaced by a 10% reflecting plane mirror aligned with the centre of the discharge. The second pulse in Figure 2.24(b) is much longer than that obtained using two PC mirrors because of an additional contribution produced by SBS of the broadband double-pass ASE radiation reflected from the real mirror at M_2 . The observation that this additional component to the output was of a much poorer beam quality than the PC component further supports this explanation since in earlier work (2.18) we found that unnarrowed KrF radiation produced a poor-quality PC reflection. Thus the presence of the second PC mirror optically isolated the system from broad-bandwidth ASE. Since Brillouin shifts are $\sim 0.2 \text{ cm}^{-1}$, and provided that the gain duration is long enough, many

cavity round trips are possible with the radiation still remaining within the $\sim 300 \text{ cm}^{-1}$ gain bandwidth of KrF.

Another application of PC mirrors that has been suggested (2.21)(2.22) is for use in the automatic alignment of targets (2.23). An example of such a scheme is shown in Figure 2.25 whereby the reflection, or glint, of the low powered oscillator from the target is sent through the amplifier and automatically retroreflected at a much higher intensity back onto the target by the PC mirror. To demonstrate this capability (see Figure 2.25) using a 50 cm lens to focus the oscillator pulse onto a 500 μm diameter plane aluminium lollipop target, which was placed 1.5 cm inside and 0.2 cm to one side of the focus of a 5 cm lens, L_1 , we were able to destroy the target with the double pass output from the amplifier by using the PC SBS mirror retroreflector. When the PC mirror was replaced by an aligned plane 80% reflector, both the reflected amplified glint and the ASE from the amplifier, which were focused by L_1 , completely missed the target and left it undamaged. Furthermore, when the target was placed at the focus of L_1 , it was destroyed by the focused broadband ASE from the amplifier. It is possible therefore that, as well as for automatic alignment of targets, this type of arrangement could be used to isolate the target from prepulse ASE arising from the final amplifier stages of a KrF fusion driver system.

2.3.3(b) Phase Conjugation of an ArF laser at 193 nm

Using an ArF laser at 193 nm we have made the shortest wavelength demonstration yet of phase conjugation. The line-narrowed oscillator/single pass amplifier system used in this study and shown in Figure 2.26(a) is similar to that used in the KrF work described above, except that the tuning element of the ArF line narrowed oscillator consisted of a 1200 line/mm grating blazed at 400 nm and was used in second order in the grazing incidence configuration. The apparent 5 cm^{-1} linewidth of the oscillator as seen in Figure 2.26(b) is instrument limited. Higher resolution measurements with a Fabry-Perot étalon showed this linewidth, $\Delta\nu_L$, to be $\sim 0.5 \text{ cm}^{-1}$. This is considerably narrower than the 10 cm^{-1} bandwidth obtained (2.24) in the only other report to our knowledge of a

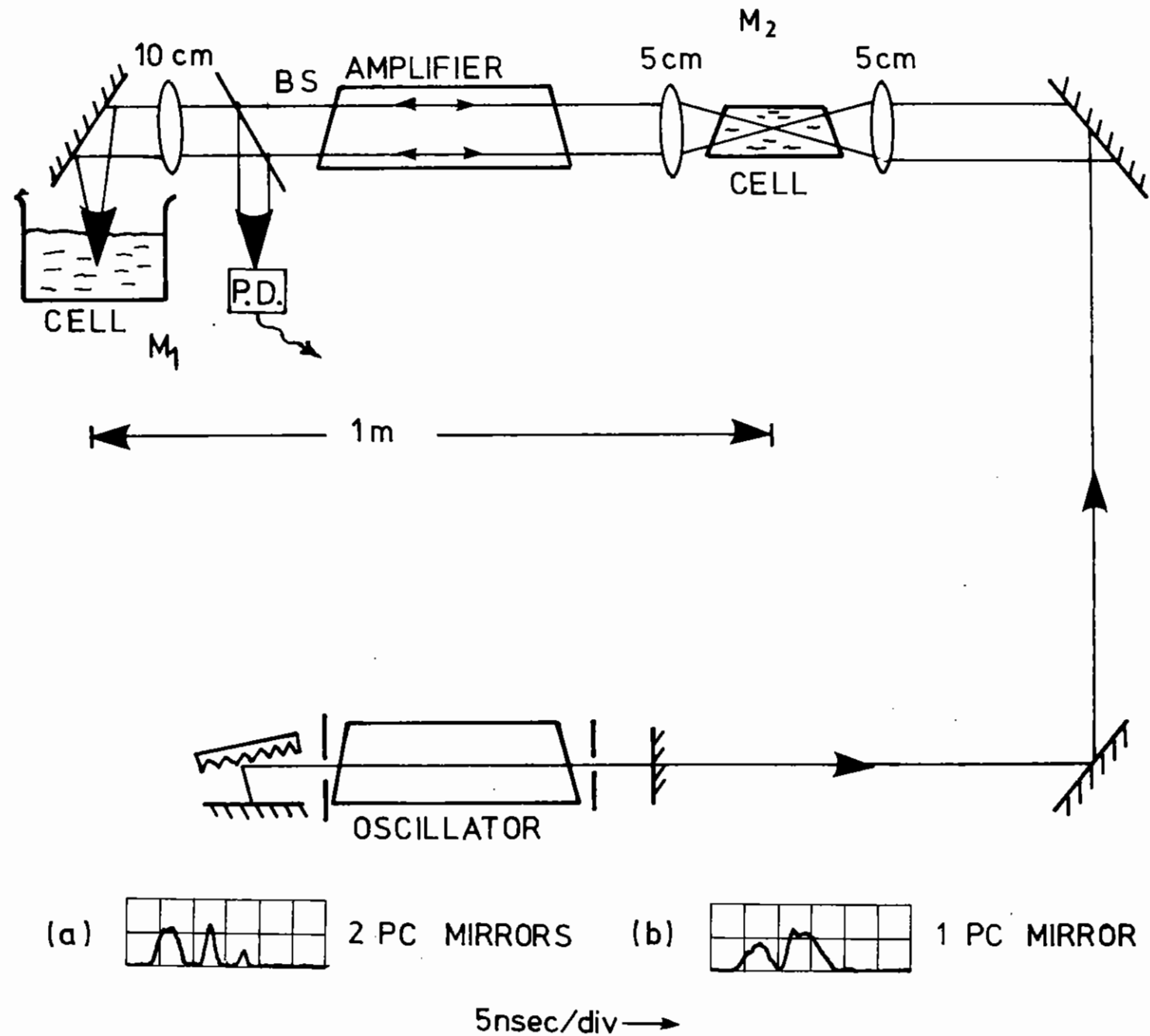


Fig. 2.24 Regenerative amplifier using SBS PC mirrors.

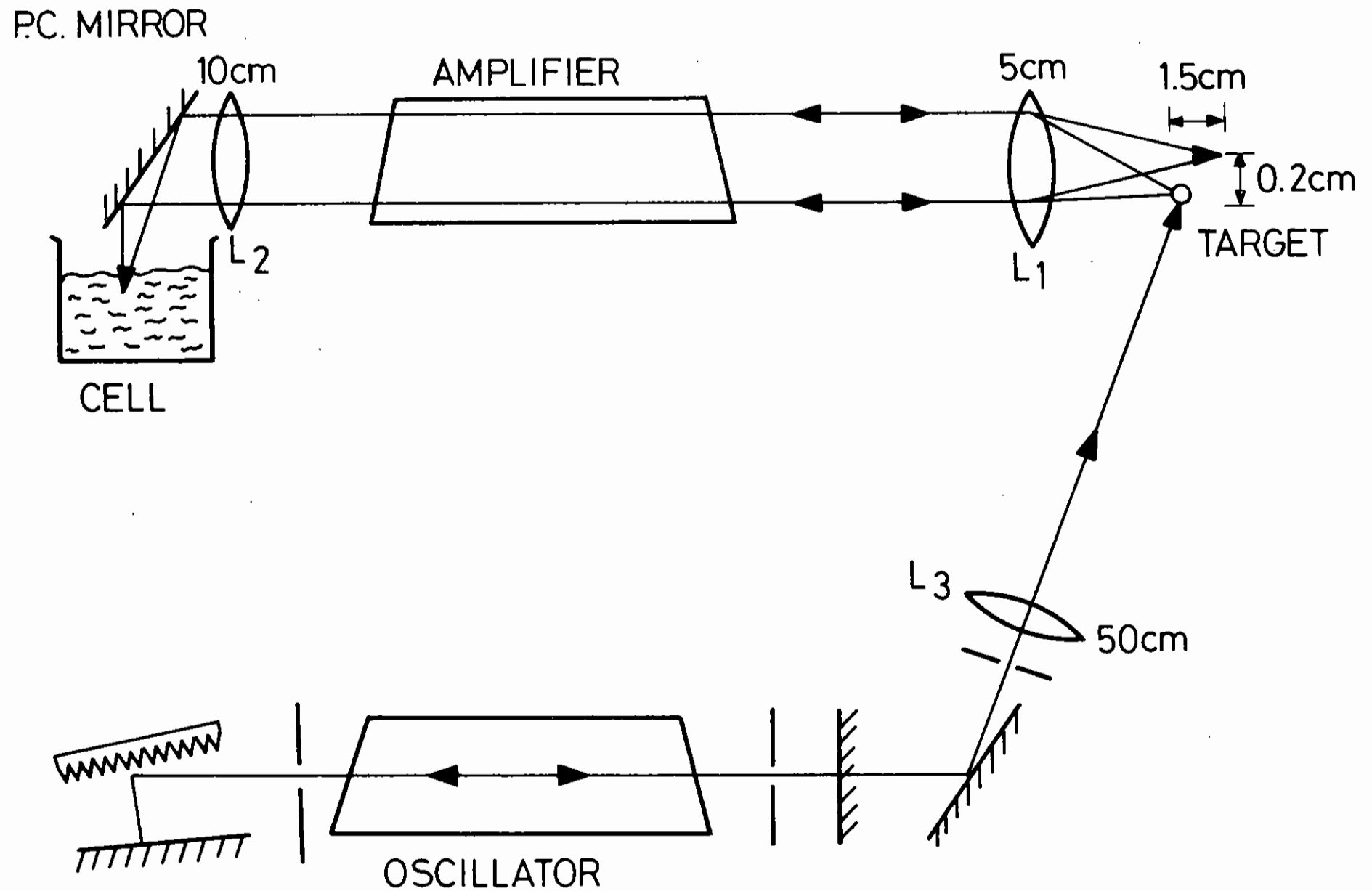


Fig. 2.25 Automatic target alignment arrangement using an SBS PC retroreflector.

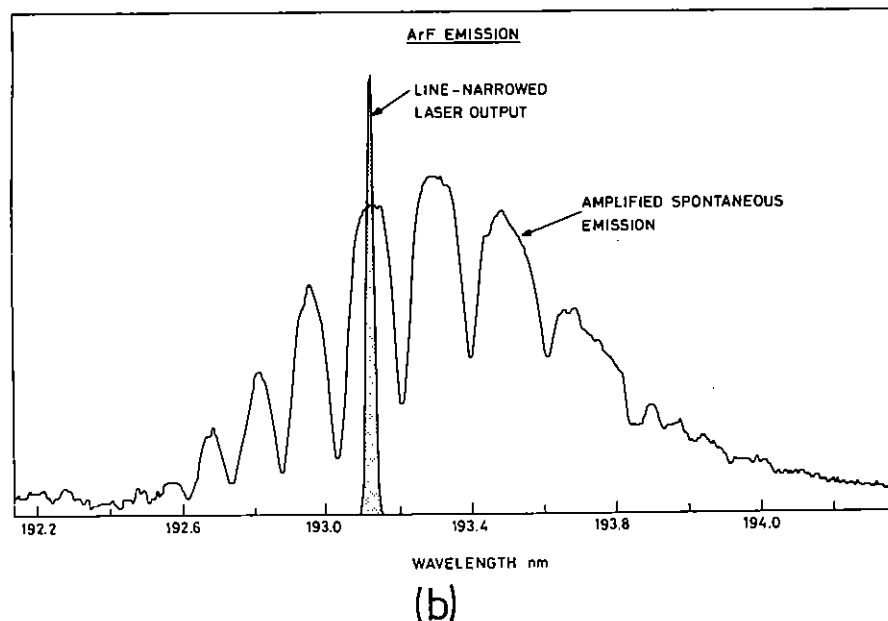
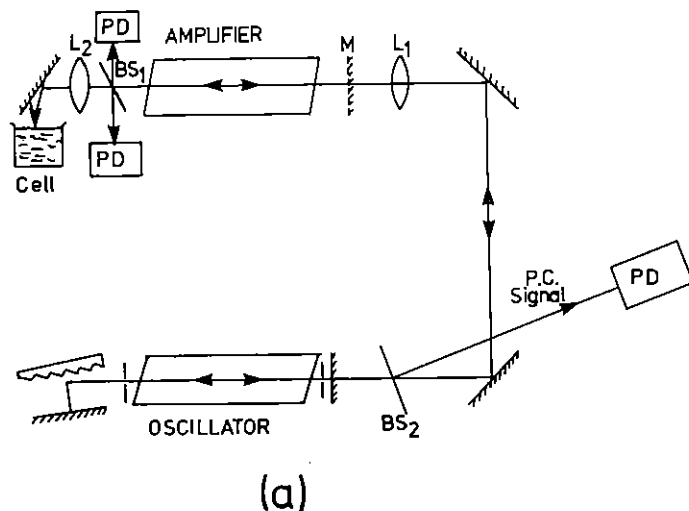


Fig. 2.26 (a) Experimental setup for ArF laser SBS PC experiments. L_1 and L_2 are 50 - and 10 - cm focal - length lenses, respectively. PD's are photodiodes.
(b) Line-narrowed output of oscillator compared with the ArF amplified spontaneous emission.

line-narrowed ArF oscillator. In that work quartz prisms were used as dispersive elements in the cavity. The line narrowed ArF oscillator output in Figure 2.26(b) is compared with the familiar ArF amplified spontaneous emission spectrum, which shows absorptions as it propagates through the atmosphere due to the (4-0) band of the $X^3\Sigma_g^- + B^3\Sigma_u^-$ Schumann-Runge system of O_2 .

To fill the amplifier volume better, the oscillator pulse (Figure 2.27(a)) is passed through a 50 cm focal length lens placed 110 cm in front of the amplifier. The resulting diverging beam passes through the amplifier with a net gain of ~ 600 . The output from the amplifier (Figure 2.27(b)) is then focused vertically by a 10 cm focal length lens into a beaker containing the scattering liquid. Beam divergences were obtained from transmitted energy measurements through pinholes of known diameter placed at the focus of long focal length lenses. To obtain absolute PC reflectivity measurements, signals were calibrated using conventional mirrors of known reflectivity substituted in place of the scattering cell.

Because the majority of commonly used solvents are highly absorbing below 200 nm, the choice of Brillouin scattering media becomes more difficult than at longer wavelengths. In the majority of the work reported here we used spectroscopic grade (BDH, Spectrosol) 2,2,2-trifluoroethanol (CF_3CH_2OH), which, after being purged with N_2 gas to reduce the absorption by dissolved oxygen, had an absorption coefficient of $< 0.3 \text{ cm}^{-1}$ at 193 nm. Furthermore, SBS was observed only when the focus of the pump radiation was $< 0.2 \text{ cm}$ below the surface of the scattering liquid. We were unable to measure the Brillouin frequency shift $\Delta\nu_B$ of the backscattered radiation with our étalon - probably because it is substantially less than our laser linewidth, $\Delta\nu_L = 0.5 \text{ cm}^{-1}$. We note that for the majority of common solvents the Brillouin shift is between 0.15 and 0.25 cm^{-1} (2.25). The wave front reversal behaviour of SBS when $\Delta\nu_L > \Delta\nu_B$ has been studied theoretically by Bel'dyugin and Zemskov (2.26). In Figure 2.27(c) we show the temporal evolution of the PC SBS pulse. It can be seen that the SBS PC mirror produces a slight shortening of the pulse. The PC pulse after it has passed through the amplifier with

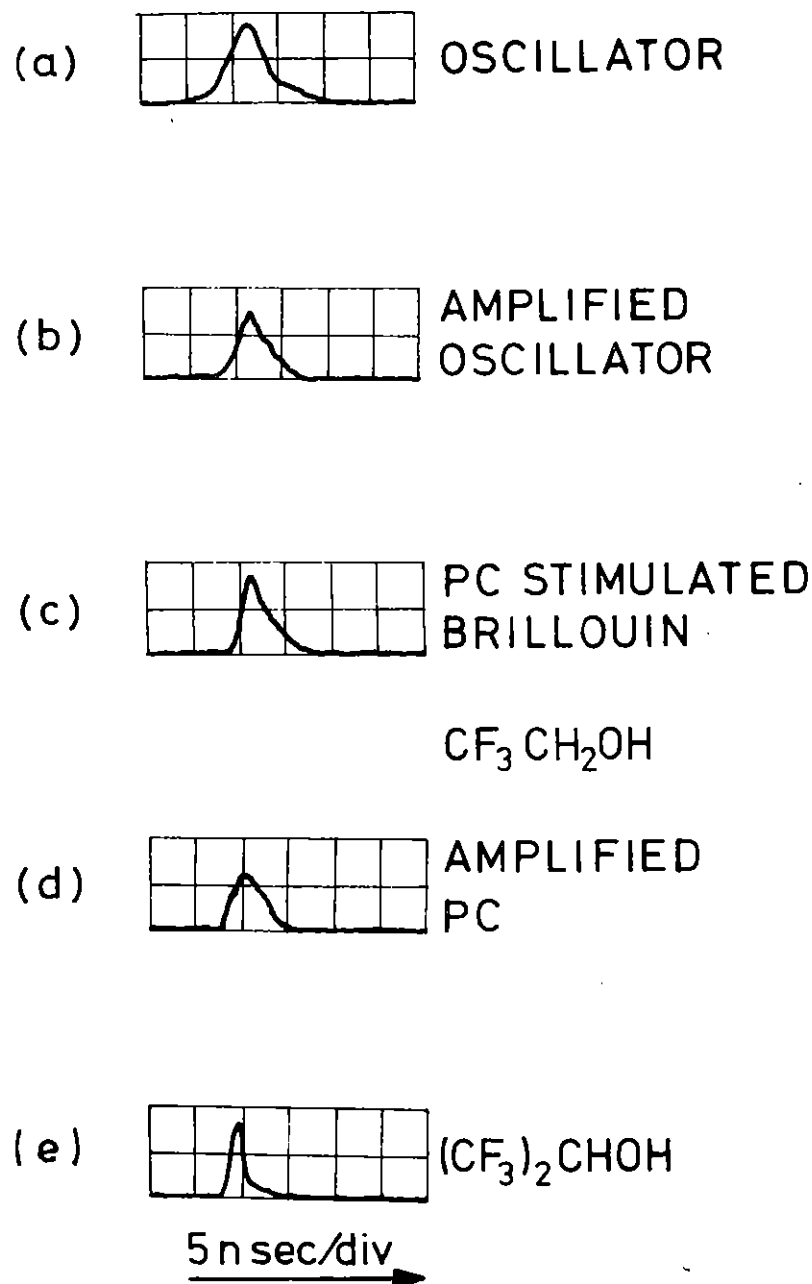


Fig. 2.27 (a) Oscillator output, (b) Amplified oscillator (SBS pump), (c) PC SBS signal from $\text{CF}_3\text{CH}_2\text{OH}$, (d) Amplified PC signal from $\text{CF}_3\text{CH}_2\text{OH}$, (e) PC SBS signal from $(\text{CF}_3)_2\text{CHOH}$.

a net gain of ~ 30 is shown in Figure 2.27(d). Whereas the beam divergence of the amplified oscillator pulse (SBS pump beam) was found to have a beam divergence of $\sim 8 \times$ the diffraction limit, from Figure 2.28 we see that the amplified PC pulse had the same nearly diffraction limited divergence as the oscillator. Thus, like the KrF system, the PC mirror corrects for the aberrations induced on the wave front by the amplifier.

The PCR of the SBS mirror is shown in Figure 2.29 as a function of the laser intensity at the focus of the pump beam in the liquid. This intensity was calculated from measurements of the beam divergence and power of the pump radiation. It can be seen that the PCR is relatively insensitive to the laser intensity - decreasing slightly at lower intensities. Although the maximum PCR measured was $\sim 11\%$ when 2,2,2-trifluoroethanol was used as the Brillouin medium, the PC mirror provides sufficient feedback into the ArF amplifier to saturate it on the return pass. Thus energy fluences and intensities of $\sim 30 \text{ mJ/cm}^2$ and 6 MW/cm^2 , respectively, were extracted from the double pass amplifier in the PC diffraction limited beam at 193 nm.

Similar PCR's were observed when 1,1,1,3,3,3-hexafluoropropan-2-ol [$(\text{CF}_3)_2\text{CHOH}$; absorption $< 0.1 \text{ cm}^{-1}$ at 193 nm] was used as the scattering medium (see Figure 2.27(e)). However, for this liquid the PCR during the latter half of the pump pulse is significantly reduced, possibly as a result of the onset of an intensity dependent absorption process, such as excited-state, photofragment, or two-photon absorption. Ethanol and hexane, which produce high quality PC SBS signals at 249 nm are highly absorbing at 193 nm and produced only poor quality and weak PC SBS returns.

M C Gower (RAL)

2.3.4 Pulse Compression and Phase Conjugation by SBS

It is well-known that stimulated Brillouin scattering (SBS) of an intense laser beam can produce phase conjugate reflection and the process has been studied both experimentally (2.27)-(2.29) and theoretically (2.30).

Recently investigations into the use of pulse compression by SBS have been undertaken (2.31)-(2.34) and the efficiency and versatility of the Brillouin mirror for this application has been demonstrated (2.34).

The characteristics of a Brillouin cavity formed by a pair of Brillouin mirrors which takes advantage of both the phase conjugating and pulse compressing mechanisms of the SBS process, has been investigated. By these means a phase conjugate resonator is produced which emits a train of intense pulses of subnanosecond duration.

The experimental investigation of the Brillouin cavity utilised the system shown schematically in Figure 2.30. The ruby oscillator was operated on a single longitudinal and a single transverse TEM_{00} mode and its output pulse was amplified by double-passing the ruby amplifier. The amplified pulses had duration ~ 40 ns, energies of < 100 mJ and linewidth of < 0.01 cm^{-1} .

The radiation was focused (at focus 1) by a 15 cm focal length lens into a chamber ($l = 157$ cm) containing methane at a pressure of 80 atm. A second short focal length lens ($f = 5$ cm) was positioned near the exit of the methane chamber to form a second focus (focus 2). The windows of the chamber were off-set at approximately 10° in order to minimise spurious reflections. By the formation of two focal regions at either end of the methane chamber, acting as Brillouin mirrors, a cavity was produced which was capable of oscillation. This will be referred to as a "Brillouin cavity".

The radiation reflected from the Brillouin cavity was monitored as shown in Figure 2.30. Figures 2.31 and 2.32 show the temporal evolution of the reflected radiation for an input laser pulse of ~ 40 ns duration and depict some of the typical features exhibited by this arrangement. For example, Figure 2.31 shows a broad pulse structure for approximately 50 ns duration transforming into a periodic train of short pulses with a repetition time of ~ 12 ns (similar to the round trip time of the Brillouin cavity). The pulse train persists for times > 100 ns.

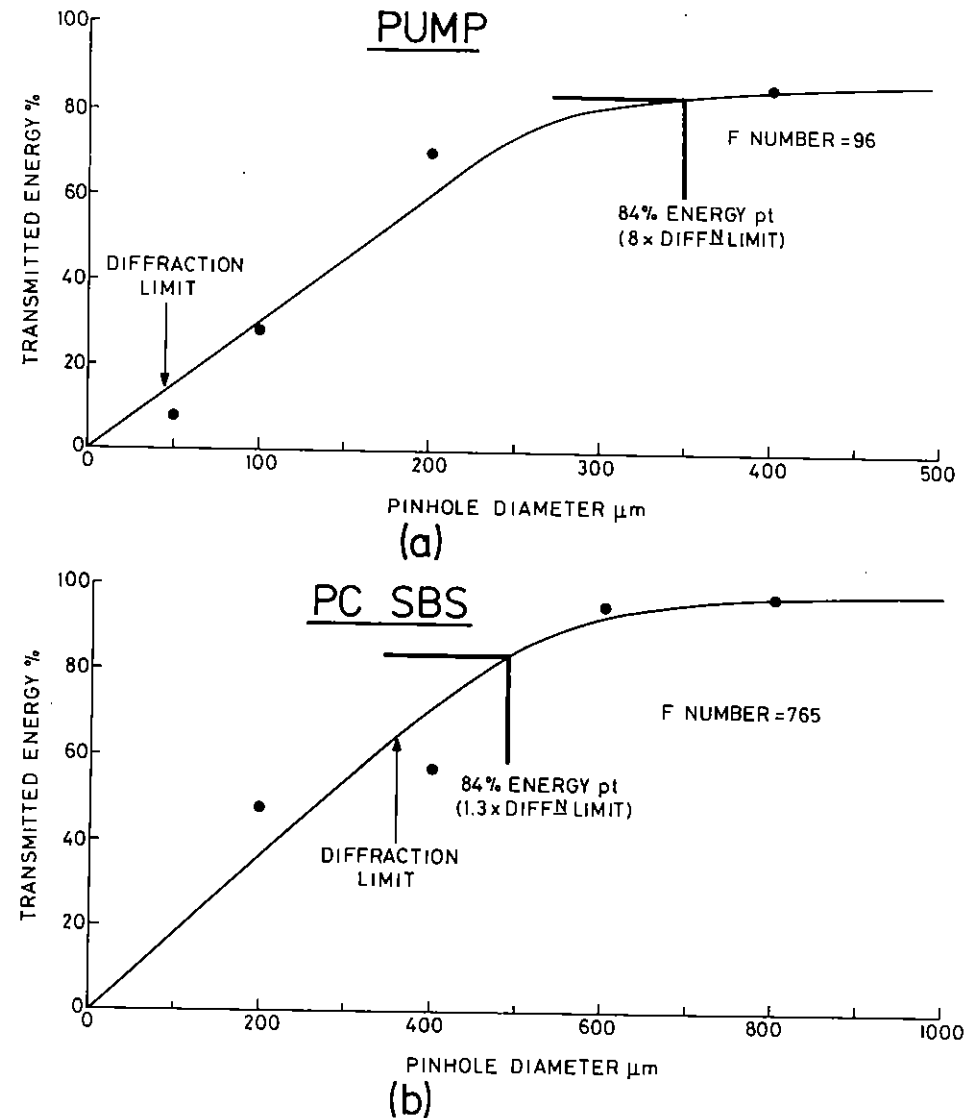


Fig. 2.28

Beam Divergences from pinhole energy-transmission measurements. The calculated diffraction-limited aperture as well as the measured 84%-energy transmission point is indicated (a) SBS pump beam. The f-number of the system is 96. (b) SBS PC beam after passage through the amplifier. The f-number is 765.

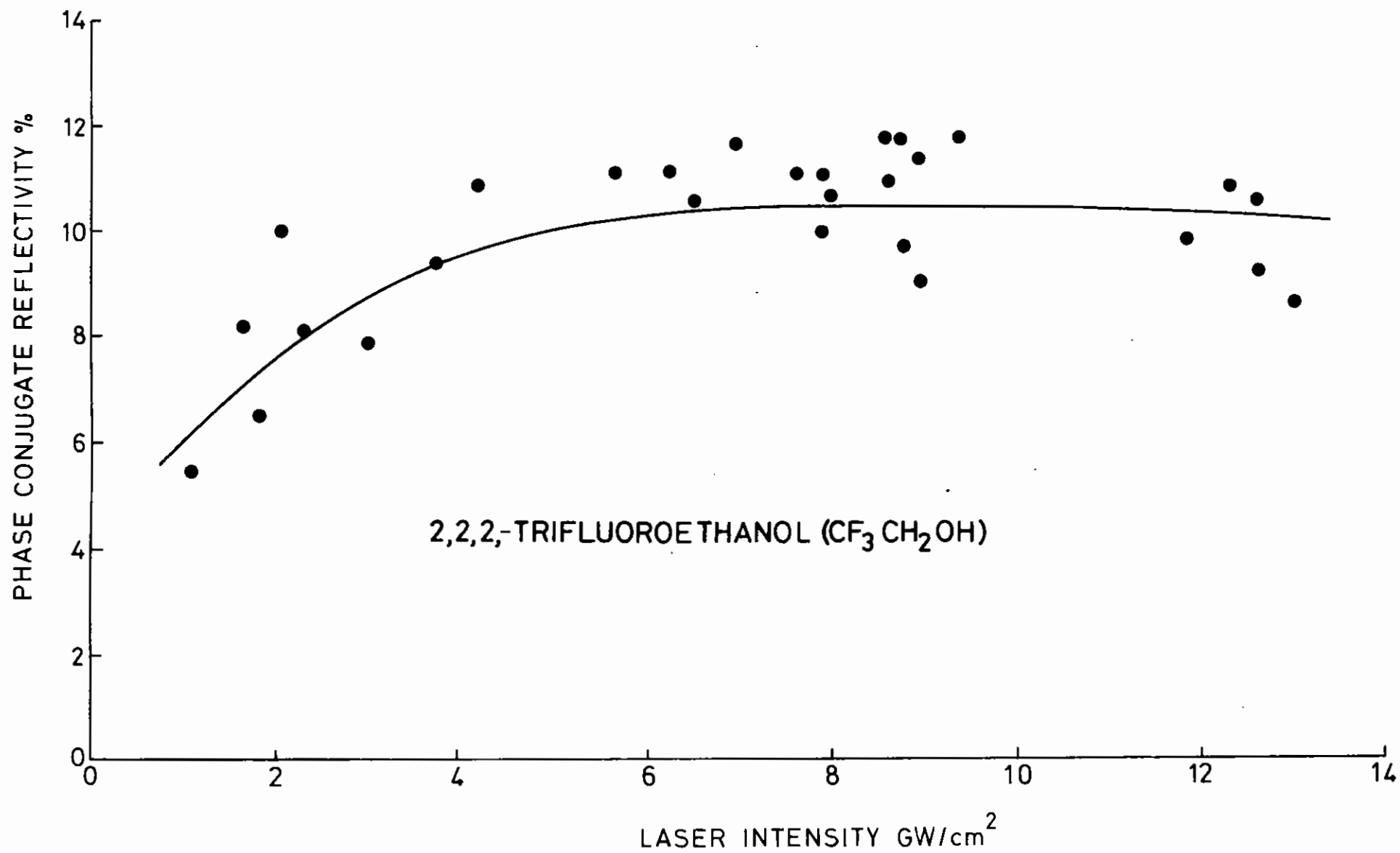


Fig. 2.29 PC reflectivity versus laser pump intensity of $\text{CF}_3\text{CH}_2\text{OH}$.

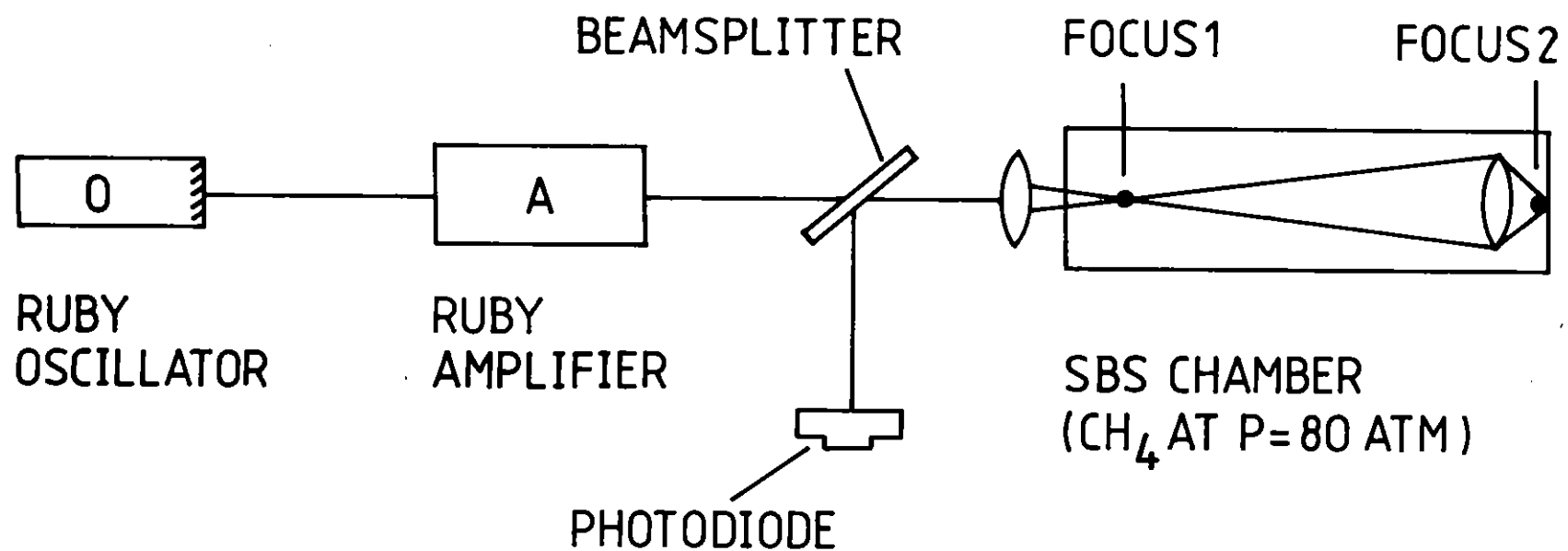


Fig. 2.30 Experimental system for investigating the Brillouin cavity.

The experimental arrangement is shown schematically in Figure 2.33(a) and Figure 2.33(b) is a simplified equivalent cavity configuration. The purpose of the ruby oscillator in the system is to inject the input narrow linewidth pulse into the arrangement to reach the threshold for SBS and initiate the oscillations. Its output mirror subsequently serves as a resonator mirror for the ruby amplifier.

The amplified pulse travelling from left to right, when focused near the front of the Brillouin cavity forms, by the Brillouin process, an acoustic wave also travelling from left to right with an acoustic intensity profile in the focal region forming a Brillouin mirror. This can backscatter the laser pump pulse with high reflectivity and the reflected pulse may also be highly phase conjugated by tight focusing (2.35). We shall for simplicity call the reflectivity of the pump pulse from the front focus, R_p , although it should be noted that the reflectivity is a function of various parameters such as pump intensity, pulse shape, focal distribution etc. Radiation reflected back from the Brillouin cavity will be reamplified in the ruby amplifier rod and be reflected by the oscillator mirror (R_0) to return to the Brillouin cavity after a round-trip time τ_p . A fraction of the pump pulse will also be transmitted through the front focal region (focus) and may then be reflected by focus 2 with reflectivity R_2 . This backscattered radiation will be refocused at focus 1 and reflected again with reflectivity R_1 . Two coupled cavities are thereby formed; a Brillouin cavity between mirrors R_1 and R_2 and a pump cavity between mirrors R_p and R_0 .

In the normal backward phase-matched SBS process, the pump wave will only interact with a copropagating acoustic wave and not with a counterpropagating one. It is therefore emphasized that R_p and R_1 although occupying the same focal region are independent reflectors under the normal Brillouin process since they are formed by oppositely travelling acoustic waves. However, the situation may be more complex than this simple model would indicate since stimulated anti-Stokes scattering can occur between counterpropagating optical and acoustic waves. In this case an acoustic photon (ω_B) and a photon (ω_0) are annihilated and a frequency up-shifted photon is created resulting in

backscattering at frequency $\omega_0 + \omega_B$. By this means a laser pulse travelling in one direction may couple to the acoustic wave deposited by a counterpropagating pulse and provide a coupling mechanism between the two resonators. Coupling may also occur in the region of focus 1 by Brillouin resonance enhanced four-wave mixing, whereby, for example, two counterpropagating waves of frequency can couple energy into a Stokes shifted wave ($\omega - \omega_B$) and an anti-Stokes wave ($\omega + \omega_B$).

The generation of the short pulses relies on several mechanisms. The nonlinear dependence of the Brillouin reflectivity on the input intensity produces pulse shortening (2.35). However pulse compression is produced by the interaction of the counterpropagating pulses in the tapered regions produced by the focusing lenses and it is likely that this mechanism is responsible for the very rapid development of the train of the short pulses observed. The pulses become progressively shorter in duration along the pulse train with subnanosecond pulses being produced under optimal conditions.

Questions arise as to the detailed mechanism whereby radiation is coupled from the regenerative ruby amplifier system into the Brillouin cavity. At the power used in the experiment, the reflectivities of the Brillouin mirrors are known to be close to unity and a pulse of radiation may be trapped within the Brillouin cavity, experiencing compression and consequent increase in power at each reflection. However it is also possible to continue to supply energy from the external laser system through focus 1. The pump radiation reflected from focus 1 may be reamplified, reflected from the output mirror of the oscillator and again be focused at focus 1. This will have the effect of sustaining the acoustic wave at focus 1 which is travelling in the same direction as the pump wave and which will provide anti-Stokes reflection for radiation within the Brillouin cavity. This means that the radiation reflected at focus 2 and of frequency $\omega_0 - \omega_B$ will experience anti-Stokes scattering with a resultant frequency $(\omega_0 - \omega_B) + \omega_B = \omega_0$. Thus the reflectivity at focus 1 can be externally enhanced and an alternate down-shift and up-shift in frequency may be achieved for some fraction of the cavity radiation.

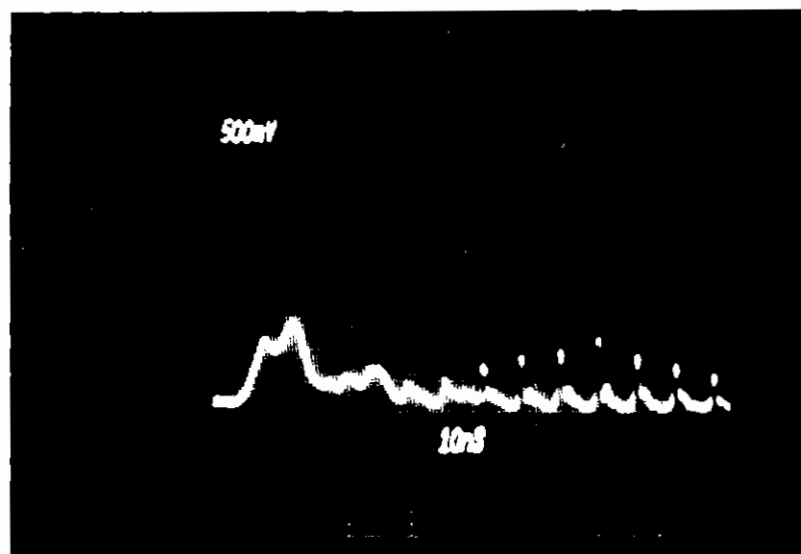


Fig. 2.31 Photograph showing backscattered pulses from Brillouin cavity.

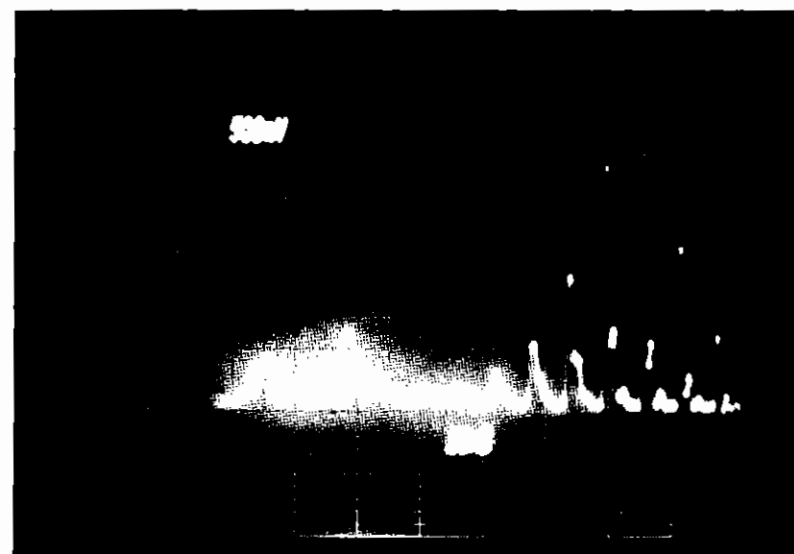


Fig. 2.32 Photograph displaying train of intense pulses emitted by Brillouin cavity.

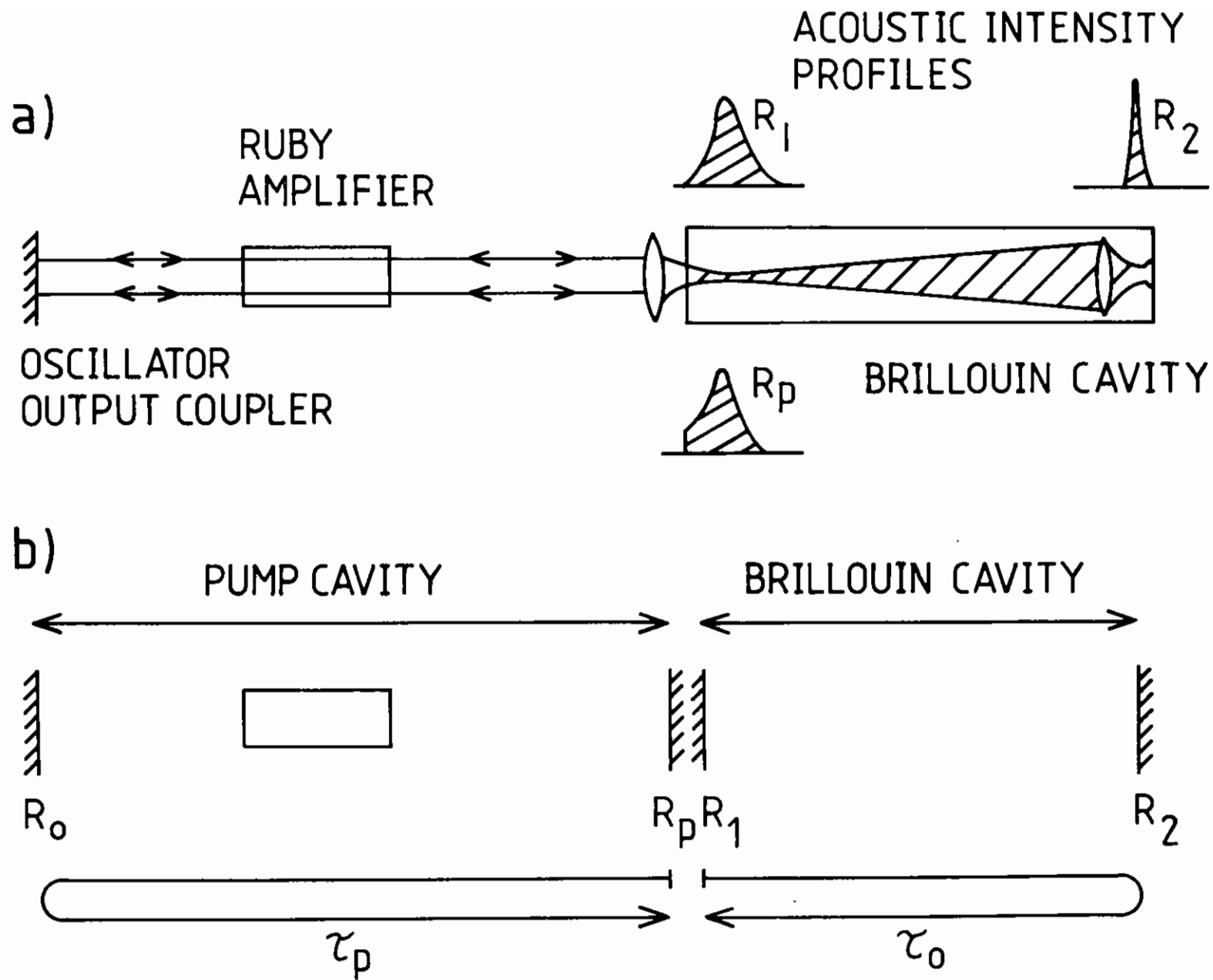


Fig. 2.33 a) Diagram depicting Brillouin cavity and coupling to laser system and b) a simplified equivalent cavity configuration.

Further aspects and applications of the Brillouin cavity are currently under study and include the synchronisation of the length of the pump cavity to that of the Brillouin cavity, an investigation of the interactions within focal region 1 where anti-Stokes and four-wave mixing interactions can occur, and the development of a numerical model to aid in the understanding of the processes.

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CHAPTER 3 UV RADIATION FACILITY

INDEX

- 3.1 INTRODUCTION page 3.1
- 3.2 DESCRIPTION OF THE UV RADIATION FACILITY page 3.2
- 3.3 PHOTOCHEMICAL AND SPECTROSCOPIC STUDIES USING HIGH POWER UV LASER RADIATION page 3.3
 - 3.3.1 Photochemical Studies
 - 3.3.2 Spectroscopic Studies of Br₂ Using VUV Laser Excitation
- 3.4 PHOTOCHEMISTRY AND RAMAN SCPECTROSCOPY page 3.8
 - 3.4.1 Introduction and Aims of the Work
 - 3.4.2 Experimental Methods
 - 3.4.3 Results and Discussion
- 3.5 PULSED UV IRRADIATION OF ZnS FILMS page 3.14
 - 3.5.1 Experimental Details
 - 3.5.2 Results
 - 3.5.3 Discussion
- 3.6 EXCIMER LASER LITHOGRAPHY page 3.19
 - 3.6.1 Introduction
 - 3.6.2 Experiments
 - 3.6.3 Results
- 3.7 UV THOMSON SCATTERING ON NON-DEBYE DISCHARGE PLASMAS AND ASSOCIATED LINESHAPE STUDIES page 3.22

- 3.8 TRIAL EXPERIMENTS page 3.23

- 3.8.1 UV Saturable Absorbers
- 3.8.2 Laser Induced Ionisation in Proportional Counters
- 3.8.3 Fluorescence of Liquid Phase Toluene Excited by a KrF Laser

REFERENCES page 3.28

CHAPTER EDITOR: F O'Neill

CHAPTER 3 UV RADIATION FACILITY

3.1 Introduction

The U.V. Radiation Facility is a new activity of the CLF which formally began operations in April 1982. The original aim was to give University researchers access to the most up-to-date U.V. laser technology for use in their research programs. In response to initial requirements the first laser installed in the facility was a commercial, discharge-excited, rare gas halide laser (Lambda Physik EMG 201). As an example of this laser's capabilities it can produce 1J laser pulses at a repetition rate of 25Hz on the KrF line at 249nm. The laser can also operate on XeF at 351 (0.4J/pulse), ArF at 193nm (0.5J/pulse) and F₂ at 157nm (15 mJ/pulse). Other major equipment available at the start-up of the facility included an optical multichannel analyzer and a Spex Triplemate Raman spectrometer.

Time on the facility is allocated to users by the UVRF Scheduling Committee - authorised users are themselves the voting members of this committee. The year is divided up into two scheduling periods, Jan. - June and July - December. In this reporting year a scheduling meeting was held on 11th May 1982 (the first such meeting) to schedule time on the facility for the period July - December 1982. A second meeting was held on the 5th November 1982 to schedule time for the period January - June 1983. At both meetings the bids for time on the facility greatly exceeded the time available but with some compromise all around, schedules were finally agreed. One $\frac{1}{2}$ day per week is set aside for general maintenance work. More lengthy periods of maintenance and facility development have to be bid for on the same basis as user experiments. Some time is also set aside as "Division Head's Weeks." This time, typically 1-3 weeks/6 month period, can be allocated to users solely at the discretion of the Head of RAL's Laser Division. This arrangement is seen mainly as a means whereby potential new users can gain access to the facility to do preliminary experiments before making more formal proposals to SERC to use the UVRF.

During the reporting year the range of equipment available at the UVRF was

gradually expanded. A Lambda Physik EMG101E discharge laser was installed for use as a line narrowed ($<0.3\text{cm}^{-1}$) oscillator. The beam from this oscillator is amplified in the EMG201 laser. An XeCl laser-pumped dye laser has also been installed for use as a probe laser in conjunction with the line narrowed UV system.

The UVRF has so far been used for various experiments in physics and chemistry with the major amount of time being devoted to photochemistry and Raman spectroscopy. As can be seen from the reports below all applications have been very productive.

F.O'Neill.(RAL).

3.2 Description of the UV Radiation Facility.

During the past year we have carried out several development which have considerably enhanced the quality of the output from the lasers in the UV facility. A schematic of the overall arrangement of the facility is shown in Fig. 3.01. We have taken delivery of a Lambda-Physik EMG101E laser module which we have modified using a grazing incidence grating type of cavity to narrow, by a factor of over 200, the frequency output from the laser. The output from this oscillator is of a diffraction limited beam quality and so is a considerable improvement when compared to the ~500 times diffraction limited quality obtained from the conventional 'off the shelf' excimer laser as obtained from the manufacturer. However, since the output energy from this oscillator is rather low (~5μJ), we use our EMG201 laser in a double-pass amplifier configuration to increase its energy. To preserve the high quality of the beam from the oscillator, we use a phase conjugate stimulated Brillouin scattering mirror as the retroreflector in this double-pass amplifier. We have previously shown (3.01) that this type of mirror (a beaker of hexane with the beam focused in it) has several advantages over more conventional mirrors - it corrects for aberrations induced on the wavefront by the amplifier and poor quality optical components, as well as isolating the system from amplified spontaneous emission (ASE).

Ultraviolet Radiation Facility

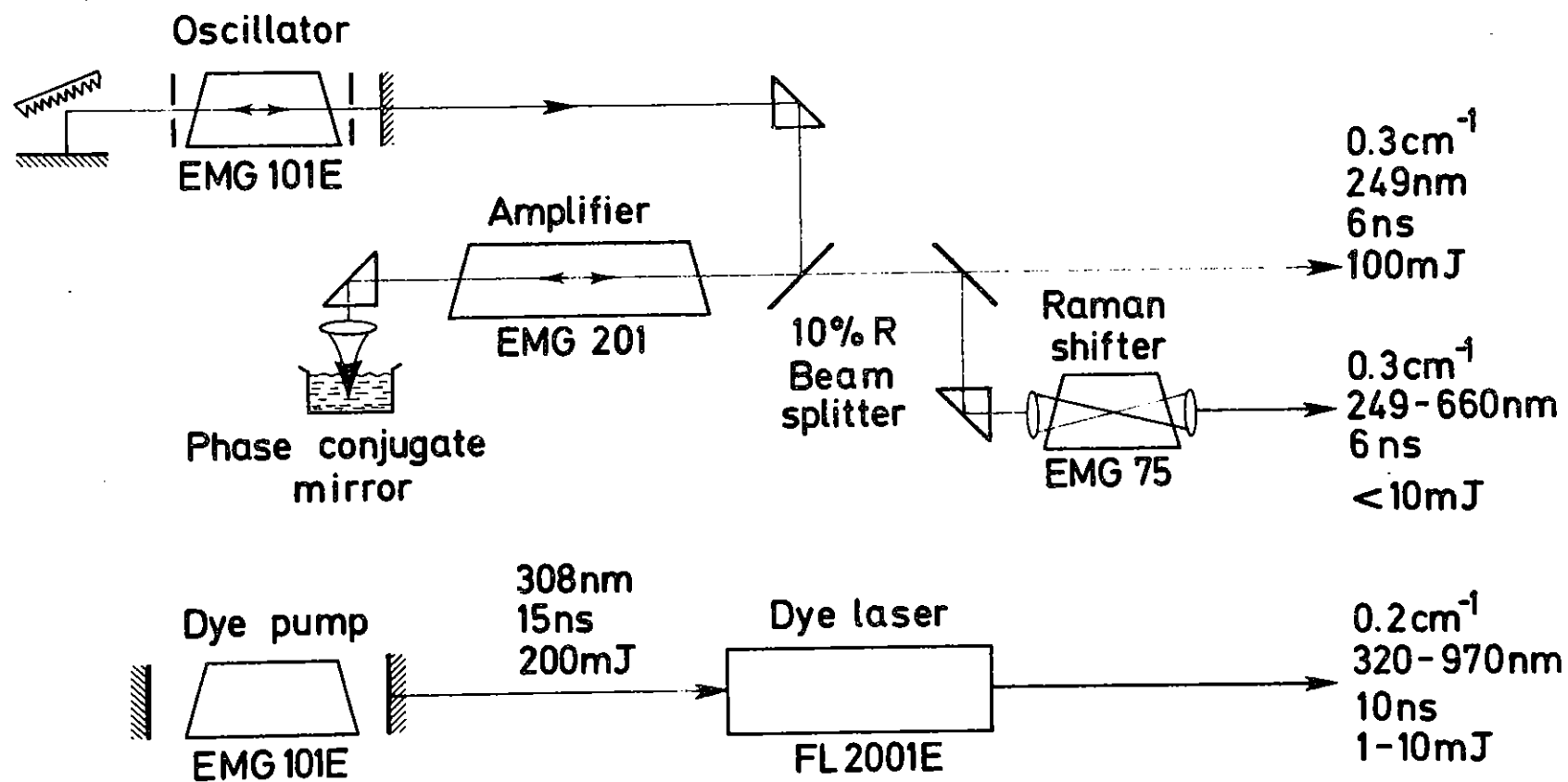


Fig 3.01 Experimental layout of U.V. Facility.

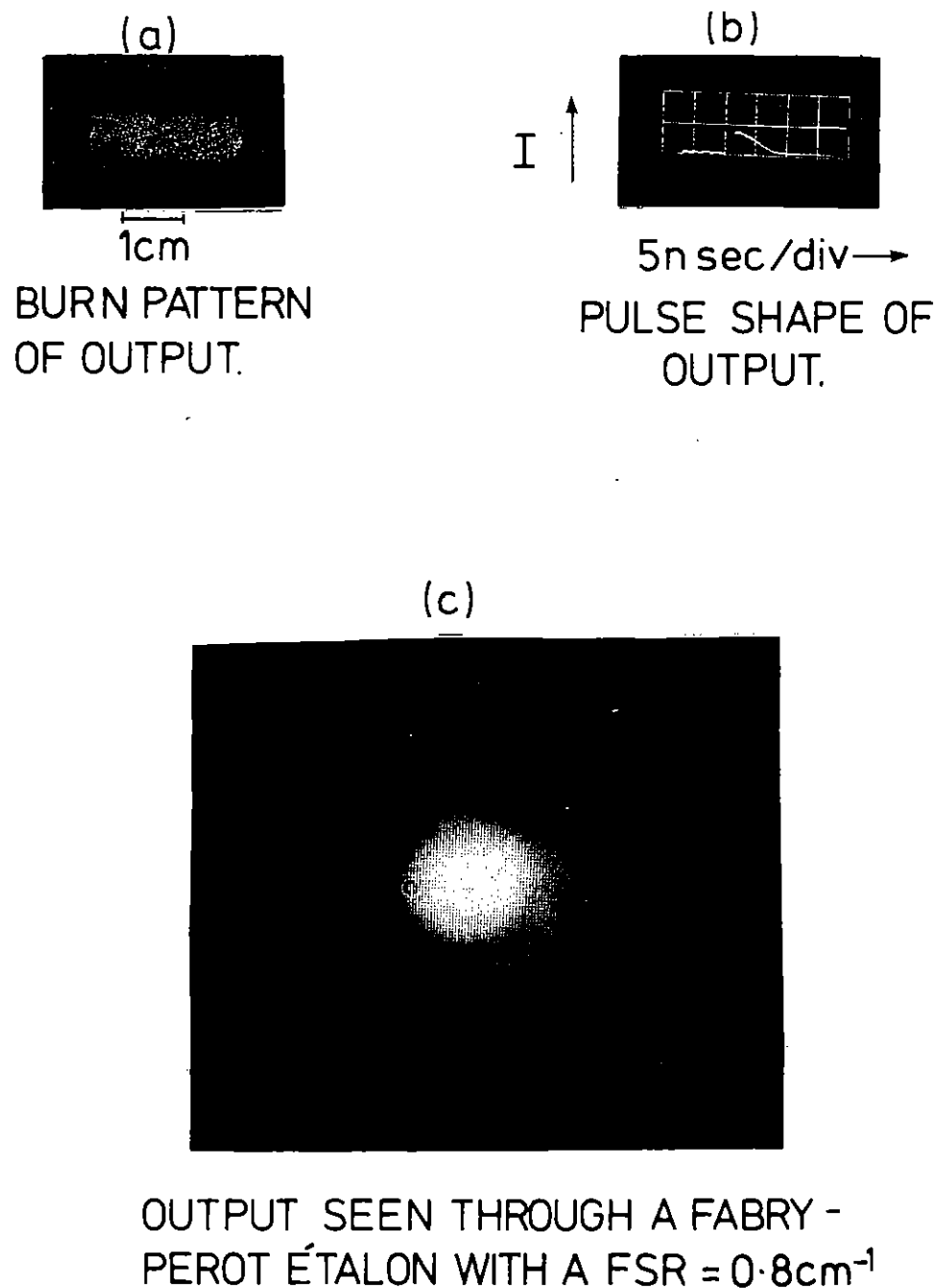


Fig. 3.02 Characteristics of line-narrowed oscillator/phase conjugate amplifier system.

Using this system we have produced output energies at 249nm (KrF) of $\sim 130\text{mJ}$ in a 6 nsec laser pulse with a near diffraction - limited beam divergence and a linewidth of $<0.3\text{cm}^{-1}$. The repetition rate can be as high as 20pps. The 'spectral brightness' (intensity per unit bandwidth) which can be achieved with this source is $\sim 10^7$ times greater than that obtainable from conventional commercial excimer lenses. As such it represents one of the highest spectral brightness ultraviolet sources in the world. In Fig. 3.02 we show a burn pattern of the uniform beam output of this system as well as an oscilloscope picture of the pulse shape. To demonstrate the narrow linewidth of the source we also show the circular fringes obtained in transmission through a low finesse Fabry Perot etalon which had a free spectral range of 0.8cm^{-1} . We will carry out further work in the near future to better synchronise the timing of this oscillator/amplifier combination.

We have also installed a Lambda-Physik FL2001E dye laser which is pumped by an XeCl (308nm) EMG101E laser. This system has a wavelength tunable output between 320-980nm with linewidth of $<0.2\text{cm}^{-1}$ and energies of 2-5mJ at a repetition rate of $<40\text{pps}$ in synchronism or with a predetermined time delay relative to the main oscillator/amplifier system.

M.C.Gower (RAL).

3.3 Photochemical and Spectroscopic Studies Using High Power UV Laser Radiation.

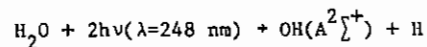
There has been considerable recent interest in non-linear spectroscopic and photochemical processes as it has been shown that previously inaccessible electronic states and novel photochemical pathways become available with non-linear techniques. In our work we are attempting to make a detailed comparison between single photon, two-photon and other non-linear processes, in order to obtain a fundamental understanding of the excitation mechanisms and primary photochemical processes involved. By studying the energy disposal in the photodissociation process it is possible to obtain information on the potential energy surface and the initial states populated in the excitation step. A broadly based review

of our work in this area, which has been carried out mainly at the Rutherford Appleton Laboratory, is given in Ref. (3.02).

The study of highly excited electronic states, whether formed by non-linear or single photon excitation, is also of interest as these can lead to new photochemical pathways. Thus in parallel with work on non-linear excitation we have carried out a series of studies using vacuum ultraviolet laser radiation. In particular we have studied the photodissociation of ICN at 157.8nm (F₂ laser) and both reactive and inelastic processes involving the ion-pair states of the halogens (I₂ and Br₂).

3.3.1(a) Two-Photon Excitation of H₂O

Two-photon excitation of H₂O at 248nm (KrF laser) leads to the formation of electronically excited OH (A²Σ⁺) radicals (3.03):



The OH(A²Σ⁺) radical can be observed in fluorescence in the region of λ ≈ 310 nm and with the high powers available from the EMG201 system the intensity is sufficient for a rotationally resolved spectrum to be obtained (see Fig. 3.02).

The populations of the individual rotational levels, for K=14-20, were readily determined using the known line strengths, as the P₁ and P₂ branches associated with these higher levels are clearly resolved. Careful analysis of the sharper features at shorter wavelength, several of which are blends of only two lines, allowed a reasonable estimate to be made of populations below K=14. These populations were then fed into a computer program which could simulate the fluorescence spectrum (3.02). Small adjustments were made to the populations until a good fit between the experimental and simulated spectra was obtained. The best fit to the data in Fig. 3.02 was achieved with a distribution that peaked at about K=16 and declined slowly up to K=20. Below K=16 the populations were

fairly constant. This distribution is clearly very different to the one observed following single-photon excitation of H₂O at 123.6 nm where a monotonically rising population of K levels, up to the maximum possible level (i.e. K=20), has been observed (3.04). The rotational distribution produced by two-photon excitation is thus somewhat cooler than that produced by single-photon excitation but the OH(A²Σ⁺) fragment still receives a considerable amount of angular momentum.

The different rotational distributions observed following single-photon and two-photon excitation suggest that different upper electronic states could be involved in the two absorption processes. In the case of single-photon excitation both the bound $\tilde{\text{C}}^1\text{B}_1$ and the repulsive $\tilde{\text{B}}^1\text{A}_1$ states are

are populated. The $\tilde{\text{C}}^1\text{B}_1$ is known to be mixed with a close lying ¹A₂ state, forming a Rydberg complex, and is also predissociated by the $\tilde{\text{B}}^1\text{A}_1$ state (3.05). Transitions to the ¹A₂ state are forbidden for single-photon absorption, but allowed for two-photon absorption. Thus two-photon excitation of H₂O can directly populate the ¹A₂ state which may therefore play a more important role in the two-photon process. As the ¹A₂ state does not correlate directly with OH(A²Σ⁺), crossing to the $\tilde{\text{B}}^1\text{A}_1$ state is still required. As this crossing must occur in a different region of the hypersurface to that for the $\tilde{\text{C}}^1\text{B}_1$ state the dynamics of the dissociation process and the energy distribution in the products will differ from those for the single-photon excitation process, despite the similar total energies involved. We intend to examine this proposal in more detail in the future using the output of the tunable, line-narrowed, KrF laser.

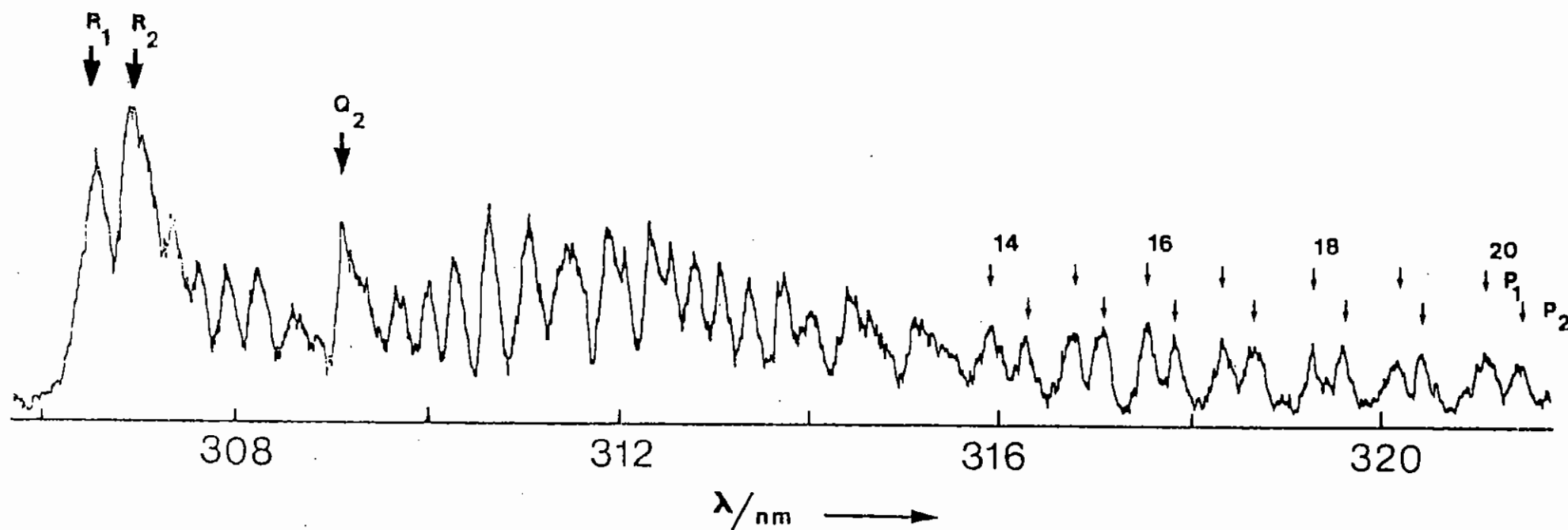


Fig 3.03 Photofragment fluorescence from $\text{OH}(A^2\Sigma^+, v'=0 \rightarrow X^2\Pi, v''=0)$ following two-photon excitation of H_2O ($P_{\text{H}_2\text{O}}=1.5 \text{ kNm}^{-2}$) with a KrF laser (248 nm). The heads of the R_1 , R_2 and Q_2 branches, together with the resolved rotational lines of the P_1 and P_2 branches, are indicated by arrows.

3.3.1(b) VUV Laser Photolysis of ICN

Relatively little work has yet been done using the vacuum ultraviolet lines available from excimer lasers. As part of a larger programme to investigate the single and multiphoton photodissociation of the cyanogen halides we have examined the photofragment fluorescence from the CN radicals formed by photolysis of ICN at 157.8 nm (F_2 laser line). Fluorescence was observed from both $CN(B^2\Sigma^+)$ and $CN(A^2\Pi_1)$, as shown in Fig. 3.04 and the vibrational energy disposal into both of these states has been determined. Our observations on $CN(B^2\Sigma^+)$ are in reasonable agreement with previous work (3.06) although we observe significantly higher populations in the vibrational levels between $v' = 3$ and 5. Fluorescence from $CN(A^2\Pi_1)$ has not previously been reported from this system and a detailed comparison of the photodissociation dynamics leading to these two different electronic states of the CN radical can now be

3.3.1(c) Reactive and inelastic Processes Involving $I_2(D^1\Sigma_u^+)$

The collisional quenching of the low lying ($E \approx 20,000 \text{ cm}^{-1}$) electronically excited states of the halogens, which are of valence character, has been studied in considerable detail. Particular emphasis has now been given to the $B^3\Pi_{0_u}^+$ state of I_2 which is readily populated from the ground state with visible radiation.

In recent work the higher lying states of I_2 , in the region of approximately $45,000 \text{ cm}^{-1}$, have also received attention. These states have a high degree of ion-pair character and their collisional behaviour may thus be expected to differ significantly from that of the lower valence states. We have recently shown that $I_2(D^1\Sigma_u^+)$ is quenched with essentially unit collision efficiency by all the noble gases and that the quenching process involves a cascade through other close lying ion-pair states. This cascade ends in what is presumed to be the lowest ion-pair state, $I_2(^3\Pi_{2g})$, and provides the basis for the operation of an

efficient optically pumped I_2 laser (3.07), which emits at 340 nm. Willard and coworkers (3.08) have shown that excitation of $I_2(D^1\Sigma_u^+)$ at 185 nm leads to efficient reaction with CH_4 and CH_3Cl , with quantum yields of 0.3 and 0.4 respectively. The high efficiencies for quenching of $I_2(D^1\Sigma_u^+)$, noted above, suggest that competition between inelastic and reactive channels could account for the quantum yields of less than unity.

The $D^1\Sigma_u^+$ state of I_2 can be conveniently pumped with ArF line at 193 nm and this populates a group of two or three vibrational levels in the region of $v' = 130$. Fluorescence from $I_2(D^1\Sigma_u^+)$ is readily monitored using an EG&G optical multichannel analyser with single shots from the laser. This has enabled us to make a wide ranging series of studies of the collisional behaviour of $I_2(D^1\Sigma_u^+)$. Of particular interest has been the observation of competition between reactive and inelastic collision processes and we have studied the related series of molecules CH_4 , CH_3Cl , CF_3Cl and CF_4 . With CF_4 by far the most dominant process is electronic energy transfer with the excited I_2 molecule being transferred through intermediate electronic states before finally reaching the lowest ion-pair state ($^3\Pi_{2g}$) from which fluorescence is the main removal process. With CH_4 and CH_3Cl reactive channels are important and lead to the formation of CH_3I and CH_2ICl . Reaction with carbon-hydrogen bonds was found to be much more rapid than with carbon-chlorine bonds, despite, the more favourable thermochemistry in the latter case. Indeed, reaction with carbon-chlorine bonds is sufficiently slow that CF_3Cl behaves essentially like CF_4 and the dominant removal process is again inter-state transfer.

The ion-pair states of the halogens clearly have interesting chemical and physical properties, and further studies are planned. A more detailed account of the above work is given in (3.09) and (3.10).

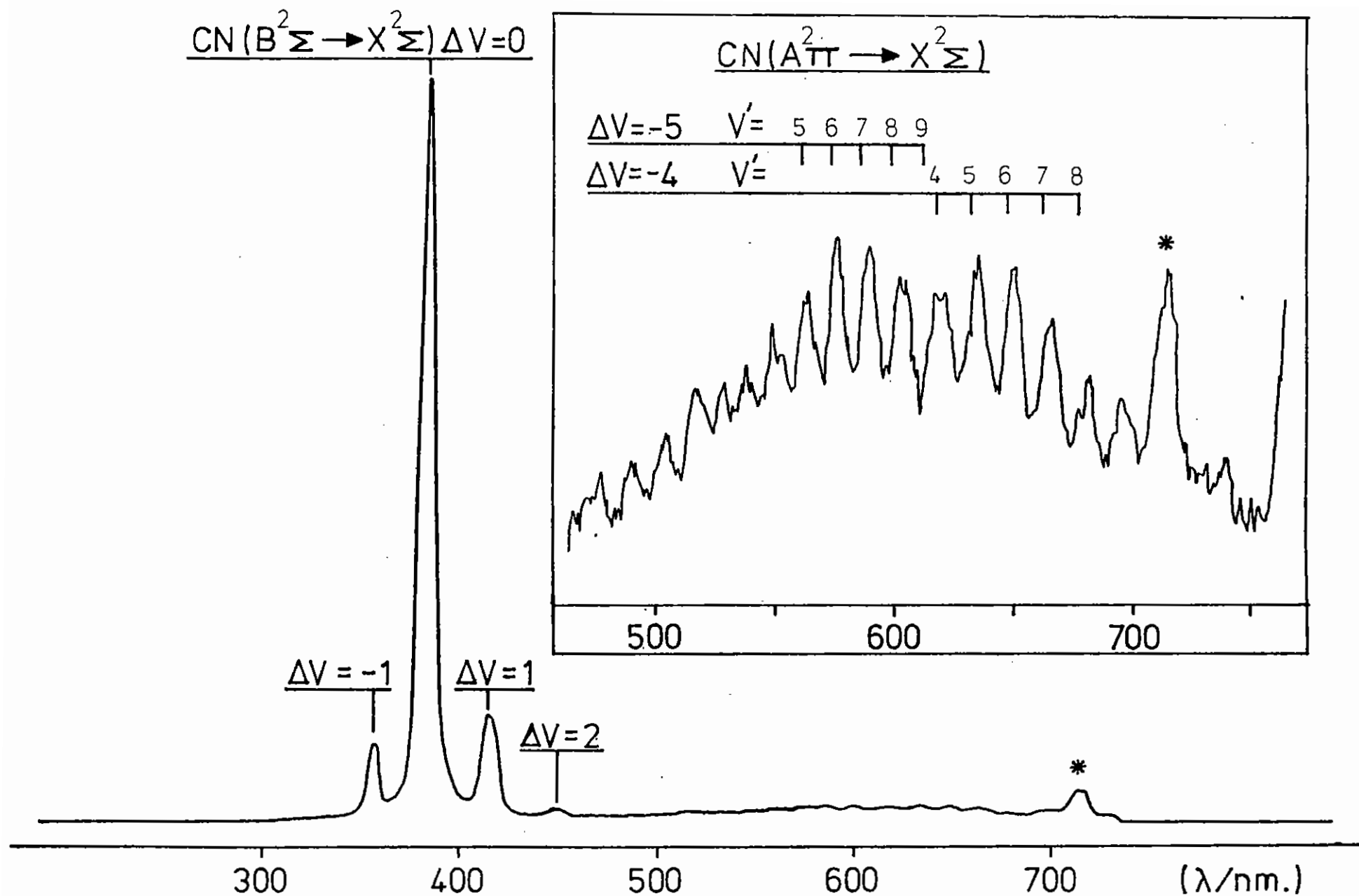


Fig. 3.04 Photofragment fluorescence from $\text{CN}(B^2\Sigma^+)$ and $\text{CN}(A^2\Pi_1)$ following photolysis of ICN at 157.8 nm (F_2 laser)

3.3.2 Spectroscopic Studies of Br₂ Using VUV Laser Excitation

Following the work described above on the ion-pair states of I₂ exploratory work on the excitation of Br₂ fluorescence, using the 157.8 nm excimer line (F₂) was initiated. Strong fluorescence, associated with three new bound-free transitions (oscillatory continua), was observed in the region 210 - 440 nm. One of these systems, extending from 300 - 360 nm is shown in Fig. 3.05. Relatively little work has yet been done on the analysis of such oscillatory continua and our work on Br₂ should provide a useful stimulus in this area. In collaboration with Dr K.P.Lawley (Edinburgh) work on the analysis of these new spectroscopic systems is in progress. Further details of the spectra and the collisional quenching behaviour of the initially populated ion-pair state of Br₂ are given in (3.11). It is clear from this work than an optically pumped Br₂ laser, similar to the very efficient system described for I₂(3.07), could be easily constructed.

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3.4 Photochemistry and Raman Spectroscopy

3.4.1 Introduction and Aims of the Work

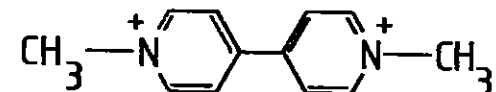
There are two general aims. The first is to exploit the line-narrowed ultraviolet output of the RAL excimer system for the excitation of Raman and resonance Raman (RR) spectra of molecular species in solution. The second is to use a controlled (in both wavelength and time) sequence of pulses from two different laser systems to initiate (by photolysis) and then to probe by (RR spectroscopy) the mechanisms of chemical reactions in solution.

Our work at the Rutherford Laboratory commenced in October, 1982, and to the time of writing this report we have had access to the UVRF for approximately 8 weeks. Although much of this time has necessarily been spent in system development and calibration we have been able to do

useful work on several different molecules and photochemical reactions. These may be summarized as follows:

3.4.1(a) Methyl viologen (MV²⁺)

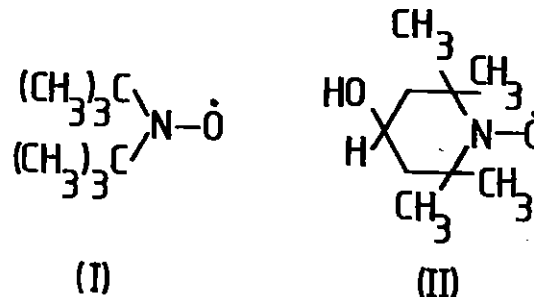
This organic compound has the structure



As the dichloride it is called paraquat and is well known for its herbicidal properties. In aqueous solution it has been much used as an electron relay in various reactions used for solar energy conversion. Its free radical cation reduction product, MV^{•+}, is capable of reducing water to hydrogen in the presence of a suitable catalyst. We have used colloidal semiconductors in the form of aqueous sols as photosensitizers for this reaction and have studied the conversion of MV²⁺ to MV^{•+} at the colloid surface by electron transfer from photopopulated conduction bands. Direct (i.e. unsensitized) photoreduction of MV²⁺ also has been observed with 248 nm irradiation.

3.4.1(b) Stable dialkylnitroxyl free radicals, R₂NO.

Ultra violet RR spectra of di-t-butylnitroxyl (I) and 4-hydroxyl -2,2,6,6-tetramethylpiperidinoxyl (II) radicals at ca. 10⁻² M concentration in mixed water/ methanol (95/5) solvent have been obtained.



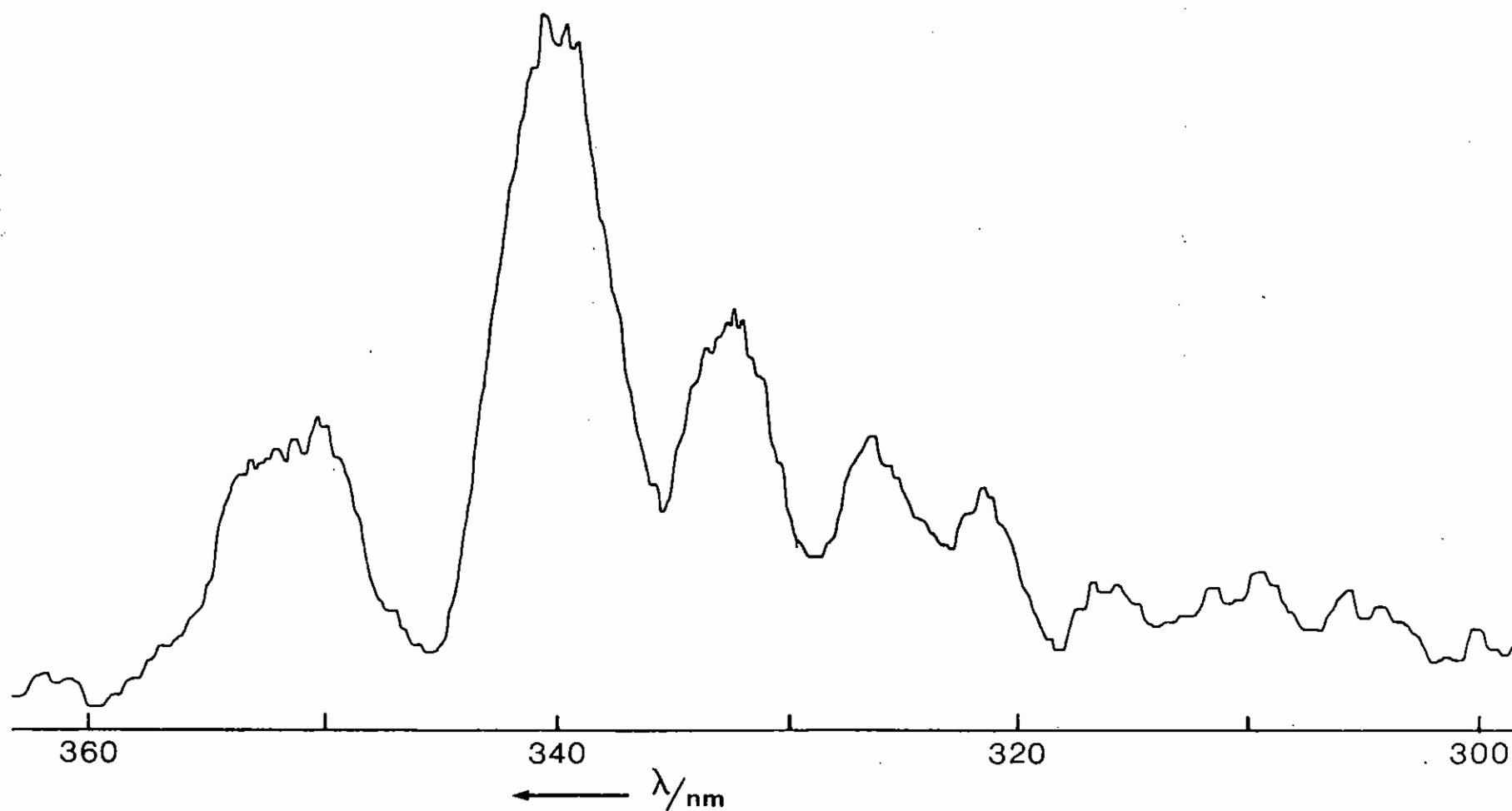
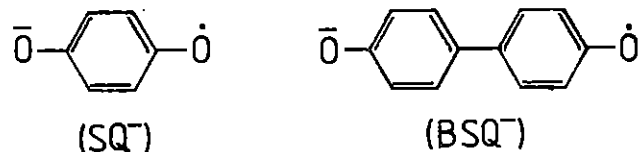


Fig. 3.05 Fluorescence from Br_2 showing one of the oscillatory continua observed by excitation at 157.8 nm (F_2 laser): $P_{\text{Br}_2} = 13.3 \text{ N m}^{-2}$.

These stable radicals are close analogs of unstable alkylperoxyl radicals which occur as important intermediates in the photo-oxidation of hydrocarbons.

3.4.1(c) p-Benzosemiquinone (SQ)⁻ and p,p'-biphenylsemiquinone (BSQ)⁻ free radical anions.

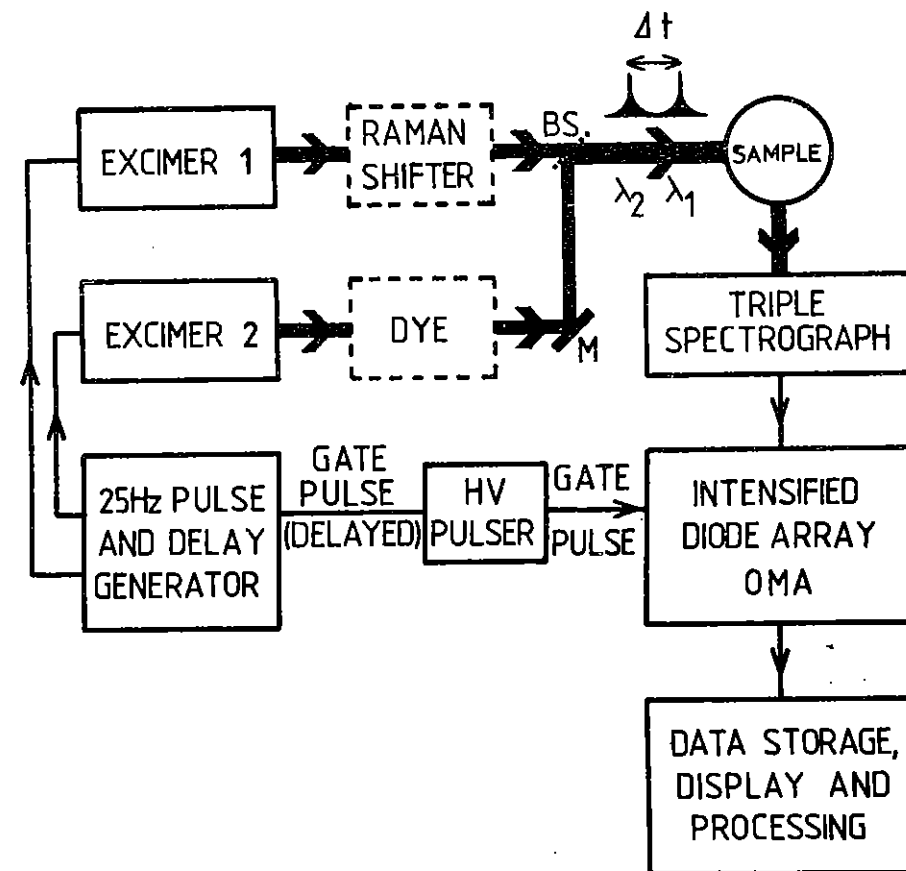


These species both have been formed as transient intermediates in the aqueous solution photolysis of their respective hydroquinone parents. The radical anions may form by a process of direct multiphoton photoionization in alkaline solution where the parent compounds are of the form ⁻O-Ar-O⁻, but their formation in acid solution depends strongly on pH and involves a chemically more complex mechanism.

3.4.2 Experimental Methods

Figure 3.06 shows in diagrammatic form the experimental arrangement of lasers, Raman spectrometer and data capture system used. This is a simplified scheme, the excimer laser component of which is specified in detail in an earlier section of this Report. The Raman shifter and HV pulser are available but have not yet been used. In addition to its function as a photochemistry 'pump' the excimer laser no. 1 also has been used in this work in its line-narrowed (to ~0.3cm⁻¹), configuration (coupled and synchronized EMG 101 and 201 KrF lasers) as a direct Raman excitation source. Since vibrational spectra of solution phase species typically contain bands of bandwidth (FWHH) ca. 10 cm⁻¹ the normal excimer laser bandwidth of ca 80cm⁻¹ is far too great to allow its use as a vibrational Raman excitation source, although it is well suited to the task of producing electronically excited states by simple absorption.

Excimer no. 2 (an EMG 101 filled with XeCl to give 308 nm laser output)



EXCIMER 1: To pump photochemical reaction in solution.

EXCIMER 2: To probe the resonance Raman spectra of short-lived reaction intermediates.

DELAY GENERATOR: To time-resolve the photochemical reaction profiles.

RAL: ULTRAVIOLET LASER FACILITY

Fig. 3.06 Diagram of equipment arrangement used for solution photochemistry and Raman spectroscopy studies.

normally is used as a pump for the FL 2001 dye laser but we also have used it directly as a photolysis pump for our samples. The line-narrowed excimer 1 then is used as the Raman probe. The dye laser output linewidth is typically ca. 0.05 cm^{-1} , which makes it immediately suitable as a Raman excitation source.

Laser pulse energies available at the sample vary with the excimer gas used but generally are sufficient to produce breakdown of the sample when focused. A few millijoules is sufficient for this. To avoid breakdown and to avoid non-linear effects our experiments usually employ partially defocused laser beams. To avoid problems of cell damage most of our work has been done with rapidly flowing jets of solution samples pumped through small syringe needles, with the laser beams impinging directly on the open jet as it emerges from the needle. Since the excimer laser pulses typically have widths (FWHH) of ca. 10 ns, only 10 mJ per pulse is equivalent to a power of 1 MW. Focused to 10 μm this gives a power density of greater than $10^{12} \text{ W cm}^{-2}$. Defocusing by several orders of magnitude is necessary to avoid stimulated Raman effects (threshold ca. 10^8 cm^{-2}).

In order to obtain good signal/noise ratios in Raman spectra excited by laser pulses of relatively low energy, repetitive pulsing of the laser is necessary. The practical limit on rep. rate is set by the EMG 201 laser (excimer 1) and this is 25 Hz. Sample flow rates of ca. 200 cm s^{-1} correspond to residence times in the 'observation zone' ('seen' by the spectrometer after magnification by the collection optics) of only ca. 20 μs , thus ensuring that a fresh portion of sample is presented to each laser pulse.

A Spex 'Triplemate' spectrometer is used for analysis of the Raman scattered radiation. This is fitted with several sets of diffraction gratings to enable it to be used over the full wavelength range from 200 to 800 nm. It consists of a double monochromator configured for subtractive dispersion, this acting as a stray light filter and bandpass selector, followed by a spectrograph stage to which is attached an intensified diode array multichannel detector (PARC model 1420 and OMA-2).

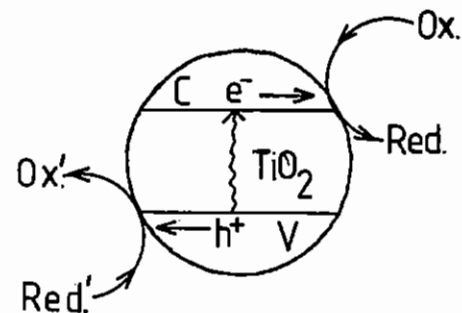
This multiplex spectrometer is well suited to the time-resolved vibrational spectroscopic work in progress. Integration of the weak Raman signals on the detector array is combined with a programmed sequence of co-additions of data in the OMA-2 processor memory. Typically ca. 4000 laser pulses (or pulse pairs in the pump-then-probe sequences) accumulated in an experiment of ca. 6 min. duration.

3.4.3 Results and Discussion

Examples of Raman spectra obtained from each of the three types of chemical system referred to the Introduction are given here. In all cases further experiments remain to be done, although the results shown are all significant in the context of the development of the methods and in providing structural and/or mechanistic information on the chemical species and reactions.

3.4.3(a) Methyl viologen, MV^{2+} , on colloidal semiconductors.

Both TiO_2 and CdS semiconductors have been used in aqueous dispersions as photosensitizers. (3.12) - (3.14). Figure 3.07 shows a resonance Raman spectrum (3.15) of the reduced species MV^+ at a concentration of ca. 1 mM adsorbed at the surface of ca. 3mM TiO_2 sol. This result was obtained by pumping the system with 248 nm (KrF excimer) radiation and, after a 100 ns delay, probing its resonance Raman spectrum with 395 nm radiation from the dye laser. 248 nm radiation is capable of promoting valence electrons through the bandgap in TiO_2 . The resulting electron-hole pairs are capable of inducing redox reactions in surface-adsorbed species according to the following scheme:



Thus it was anticipated that MV^{2+} could be reduced to MV^+ at the TiO_2 colloid surface.

A similar, though weaker, spectrum of MV^+ was obtained using 193 nm (ArF excimer) pump radiation in place of 248 nm and the spectrum showed a fall-off in intensity as the delay between the pump and probe laser pulses was increased to ca. 150 μs . This decay of the MV^+ spectral intensity could be attributed to reaction of the MV^+ radical with dissolved oxygen. However, 248 nm also coincides with an adsorption band (3.15) of MV^{2+} and we have found that even in the absence of the TiO_2 sensitizer pumping MV^{2+} solutions with the 248 nm laser also results in formation of MV^+ radical spectrum. The mechanism for this photoreduction is not known but is currently under investigation.

3.4.3(b) Nitroxyl and alkylperoxyl radicals.

Figures 3.08 and 3.09 show resonance Raman spectra of di-*t*-butylnitroxyl and 4-hydroxy-2,2,6,6-tetramethylpiperidinooxyl free radicals, respectively, each at 1.1×10^{-2} M concentration in 95% H_2O /5% CH_3OH solvent. These were excited by the line-narrowed KrF laser radiation at 248 nm which coincides with a strong adsorption band associated with the NO chromophore. Thus the NO bond stretching mode is most strongly resonance enhanced in these Raman spectra and gives the bands at 1346 and 1361 cm^{-1} , respectively. These band wavenumbers may be used to determine N-O bond force constants and to infer the substituent effects in the two radicals.

A preliminary experiment to observe the RR spectrum of an alkylperoxyl radical was unsuccessful. This involves two lasers, one at 308 nm to generate the species in situ by photolysis and the other (the probe laser) at 248 nm to coincide with an absorption band of the peroxyl radical and thus to excite its RR spectrum. However, it has been possible to identify the reasons for failure and it is hoped to overcome the problems in the near future.

3.4.3(c) p-Benzosemiquinone and p,p'-biphenylsemiquinone (BSQ⁻) free radical anions.

Figure 3.10 shows the RR spectrum of the BSQ⁻ radical formed from its parent hydroquinone at 10^{-2} M concentration in strongly alkaline aqueous solution. The photoionization reaction was pumped by the KrF 248 nm excimer laser and the RR spectrum probed with 441 nm dye laser excitation. Delays between the two laser pulses from 20 ns to 5 μs were used and the spectrum remained strong throughout, indicating a rapid formation and good stability for this reaction intermediate. The spectrum is closely similar to that produced previously by photolysis of the hydroquinone at low temperature in glassy ethanol matrices (3.16).

Similar data have been obtained from benzosemiquinone (3.16) (3.17). These enable comparisons with parent spectra to be made and conclusions to be drawn regarding changes in electronic structure and chemical bonding which accompany the photooxidation reactions. Further experiments are planned to explore the effect of changing pH on the spectra and thus on the photoreaction mechanisms.

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3-5 Pulsed UV Irradiation of ZnS Films

Most of the effort on laser annealing of ZnS:Mn thin films since the 1982 report to the LFC (Section 2.4.3) has been devoted to a study of the structure of the annealed material using reflection high energy electron diffraction (RHEED) measurements. This has been investigated for a variety of film/substrate, film deposition and annealing parameters. Quantitative dopant profiles, before and after annealing, have been obtained using secondary ion mass spectrometry (SIMS).

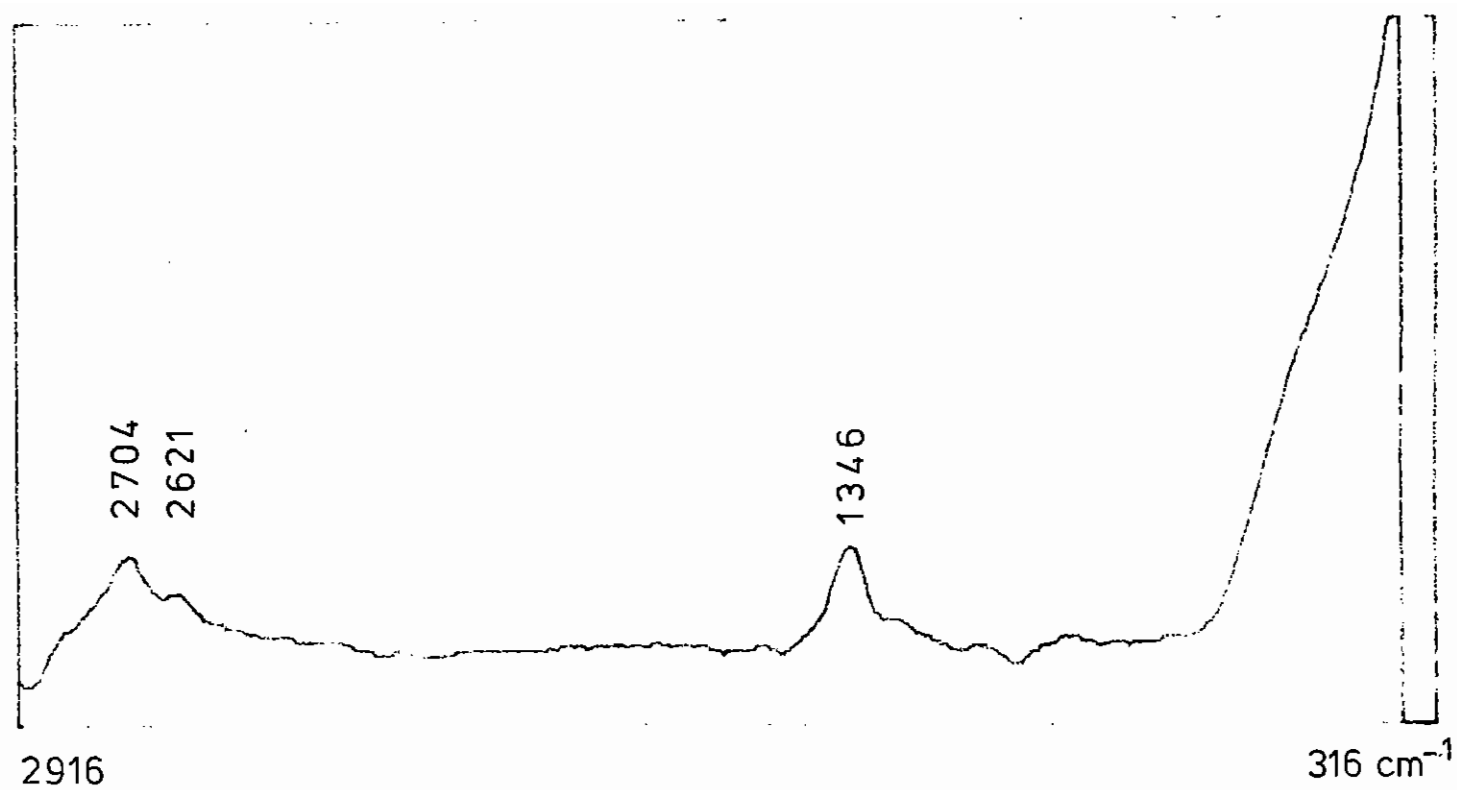


Fig. 3.08 RR spectrum of the di-*t*-butyl nitroxyl radical at 1.1×10^{-2} M in 95% H₂O/5% CH₃OH. Excitation with line-narrowed KrF laser at 248 nm.

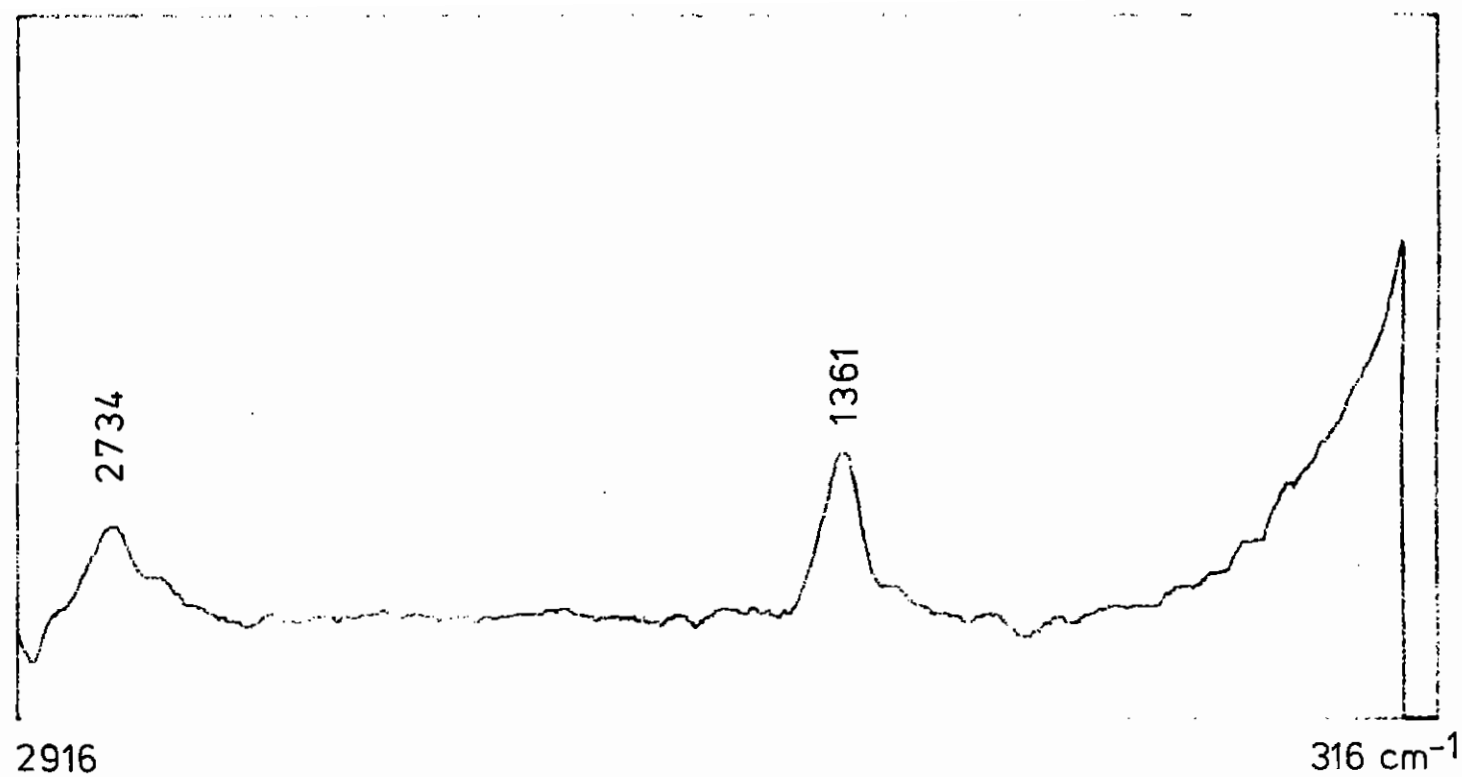


Fig. 3.09 RR spectrum of the 4-hydroxy-2,2,6,6-tetra-methylpiperidino oxyl radical at 1.1×10^{-2} M in 95% H_2O /5% CH_3OH . Excitation with line-narrowed KrF laser at 248 nm.

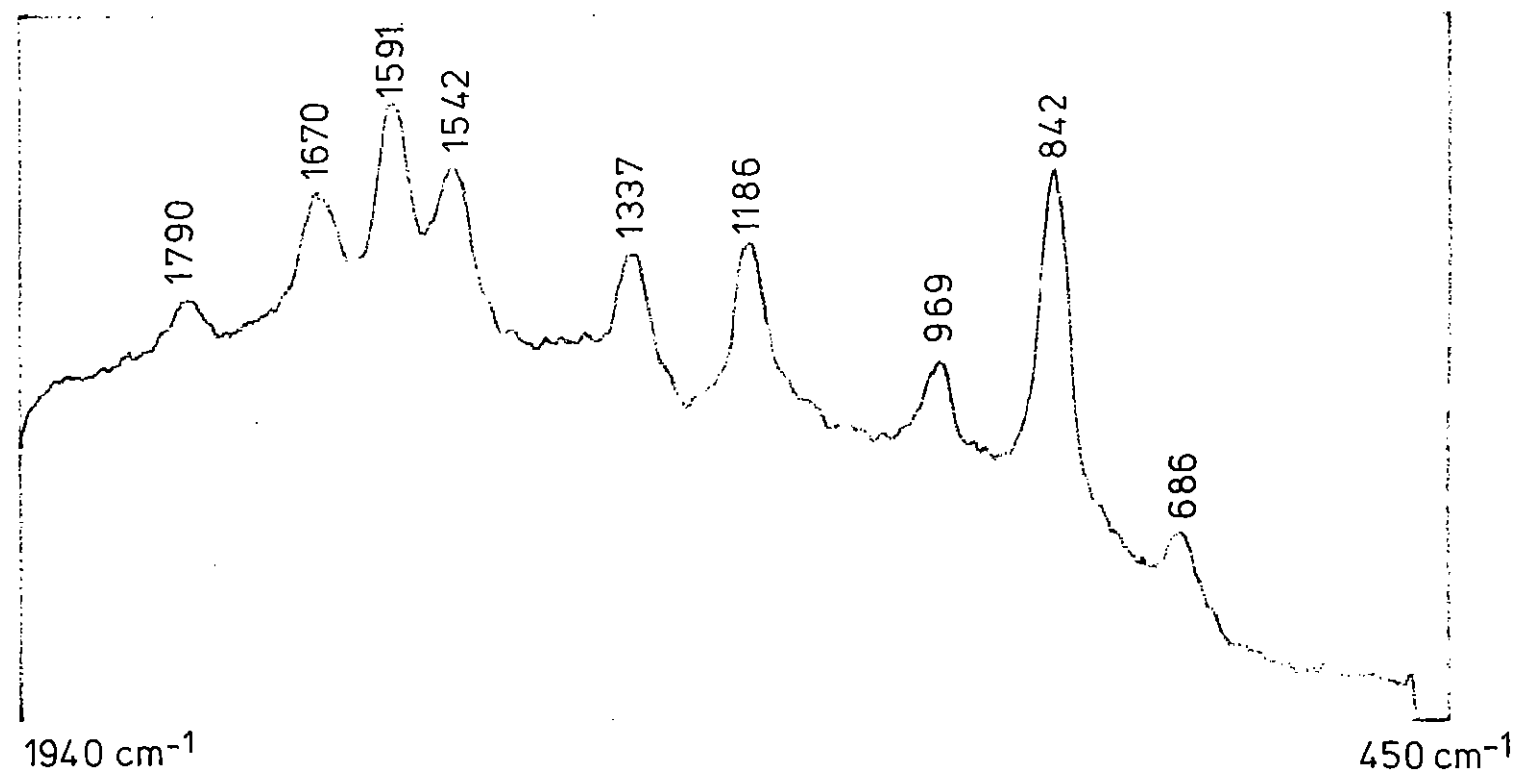


Fig. 3.10 RR spectrum of the p,p' - biphenylsemiquinone radical produced by photolysis of its hydroquinone parent at 10^{-2} M in strongly alkaline aqueous solution. 248 nm pump, 20 ns delay, 441 nm Raman probe.

3.5.1 Experimental Details

The ZnS films were deposited onto n-type (100) Si wafers either by evaporation from solid source material in an ultra-high vacuum (UHV) system or by radio-frequency (RF) sputtering from a solid target in an argon atmosphere. Although Si has a good lattice match with the cubic phase of ZnS, the mismatch in the thermal expansion coefficients is not insignificant. For this reason films were also deposited, by RF sputtering only, onto (100) semiconducting n-type GaAs wafers for comparison purposes. This material has a better thermal match with ZnS but the lattice match is poorer. The relevant thermal and lattice data is summarised in Table 3.01.

Films deposited by evaporation in UHV onto oxide-free Si substrates were single crystalline; all the sputtered specimens were polycrystalline. The film thicknesses employed ranged between $\sim 0.2 - 1.0 \mu\text{m}$. Selected specimens were implanted with 300 KeV Mn ions at $\sim 3-5 \times 10^{15} \text{ ions cm}^{-2}$ at A.E.R.E. Harwell. The laser anneals were carried out under pressure using the 'ELF' excimer laser system at RAL and the procedure described previously (3.18, 3.19). In addition to RHEED and SIMS measurements, specimen surfaces were also studied optically using Nomarski interference microscopy (3.20).

3.5.2 Results

Figure 3.11 shows characteristic RHEED patterns from $\sim 0.8 \mu\text{m}$ thick single crystal films on Si subjected to various treatments. The data for the virgin film, Figure 3.11(a) shows a spot pattern corresponding to cubic ZnS but additional reflections are observed. These can be accounted for by invoking the presence of twinning and double diffraction effects but the presence of a small amount of hexagonal material cannot be ruled out. Results obtained after Mn implantation and a 1 hour, 500°C thermal anneal of the implanted film are shown in Figure 3.11(b) and 3.11(c) respectively. From the re-emergence of the spots in Figure 3.11(c), thermal annealing appears to be effective in restoring some structural order. It is also interesting to note that there is no hint of the

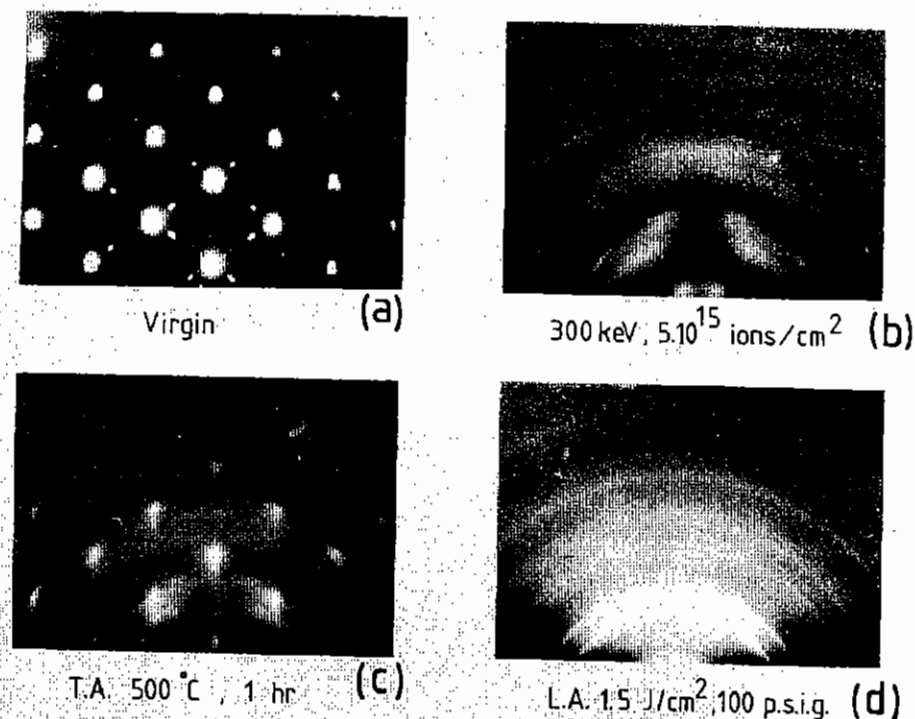


Fig. 3.11 RHEED patterns from $\sim 0.8 \mu\text{m}$ thick single crystal films on Si (a) Virgin specimen (b) After Mn implantation (c) Implanted specimen after thermal annealing (d) Implanted specimen after a 1.5 J cm^{-2} XeCl pulse under 100 psig Ar.

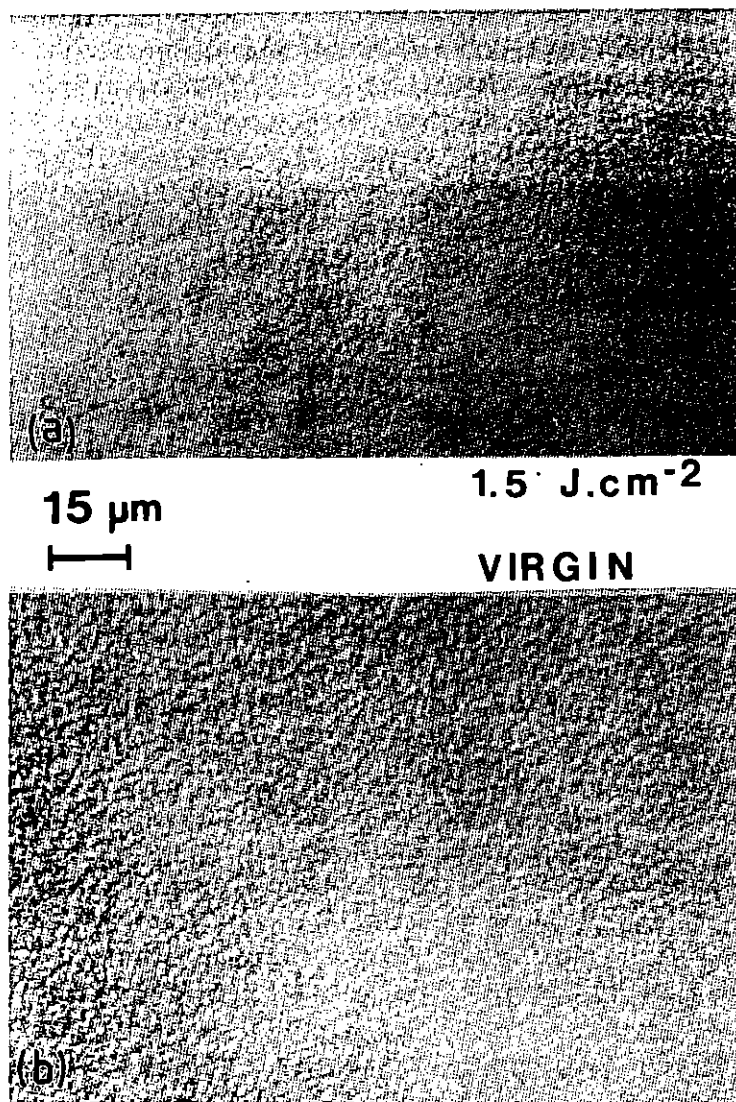


Fig. 3.12 Normarski micrographs of specimens of Figure 1: (a) Laser annealed film
(b) Virgin, as-implanted or thermally annealed films.

additional satellite spots of Figure 3.11(a) in Figure 3.11(c), the latter being characteristic of purely cubic material. The data shown in Figure 3.11(d) was obtained after the Mn implanted films had been irradiated with a 1.5 J cm^{-2} XeCl pulse under ~ 100 psig of Ar; sharp rings characteristic of randomly orientated crystallites are observed.

Nomarski micrographs taken on the samples of Figure 3.11 are shown in Figures 3.12(a) and (b). The latter shows the relatively rough surface appearance of the virgin films. As-implanted and thermally annealed specimens had an identical appearance. On the other hand, the micrograph of the laser annealed specimen shown in Figure 3.12(a) is significantly different, exhibiting a cellular structure. Behaviour similar to the above has been seen on all our epitaxial films on Si, which lay in the thickness range $\sim 0.6 - 1.0 \text{ μm}$, up to XeCl annealing energy densities of $\sim 2 \text{ J cm}^{-2}$. Thinner epitaxial specimens have not yet been investigated.

A more extensive study has been conducted on sputtered films using both XeCl and XeF radiation. For this two film thicknesses of $\sim 0.2 \text{ μm}$ and $\sim 1.0 \text{ μm}$ were employed; also both undoped and Mn-implanted films were investigated. All the As-deposited films were polycrystalline with a cubic structure and varying amounts of fibre texture. The latter was most pronounced when Si which had been argon-ion etched prior to deposition was used as the substrate. Table 3.02 summarises the RHEED results obtained after laser annealing these specimens.

In all cases polycrystalline regrowth is obtained irrespective of substrate preparation, type or film doping effects. Any initial fibre texture in the films is eliminated except for the 1 μm films where it decreases with increasing annealing energy density. Typical RHEED patterns for the 0.2 μm films on either Si or GaAs taken below and above $\sim 2 \text{ J cm}^{-2}$ are shown in Figures 3.13 (a) and (b). The former exhibits rings additional to those of cubic ZnS, like the data of Figure 3.11(d), which are consistent with the presence of either hexagonal ZnS or, possibly, double diffraction effects. Figure 3.11(b) on the other hand suggests the presence of purely cubic phase material; the spots on the rings imply a larger grain size relative to films annealed at lower energy

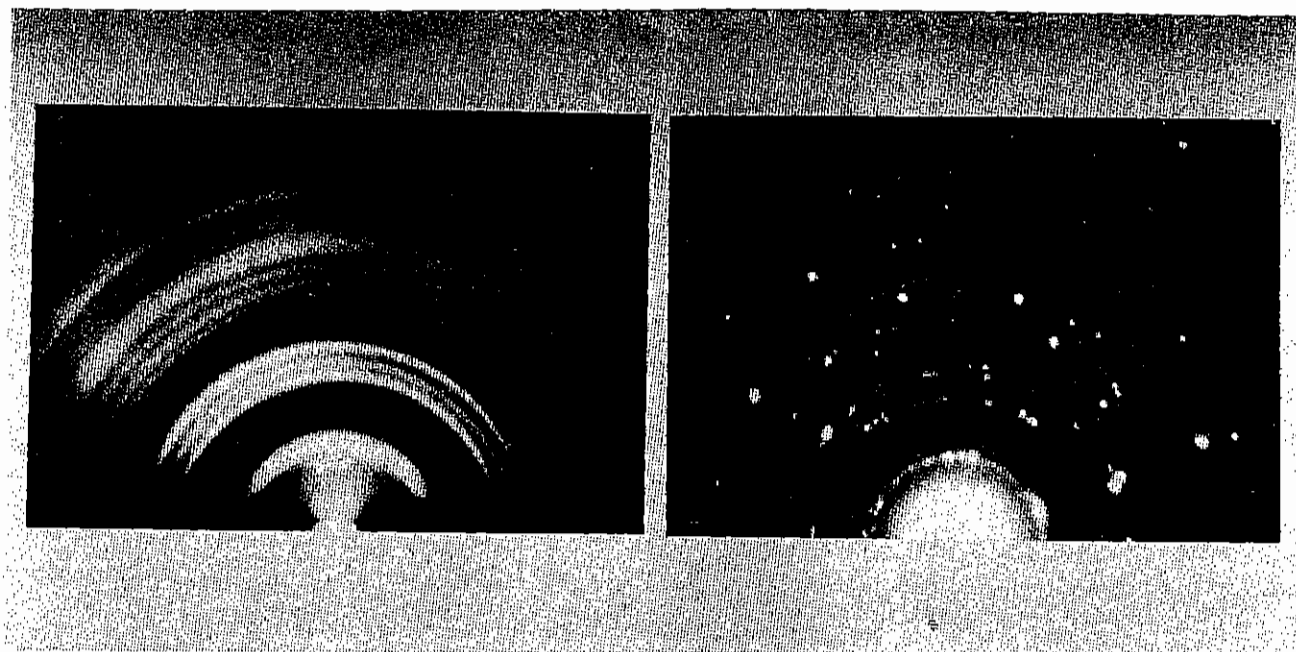


Fig. 3.13 Typical RHEED patterns from laser annealed $0.2\text{ }\mu\text{m}$ thick films on Si or GaAs (a) 1.5 Jcm^{-2} (b) 2.3 J cm^{-2} . XeCl radiation, 150 psig Ar pressure.

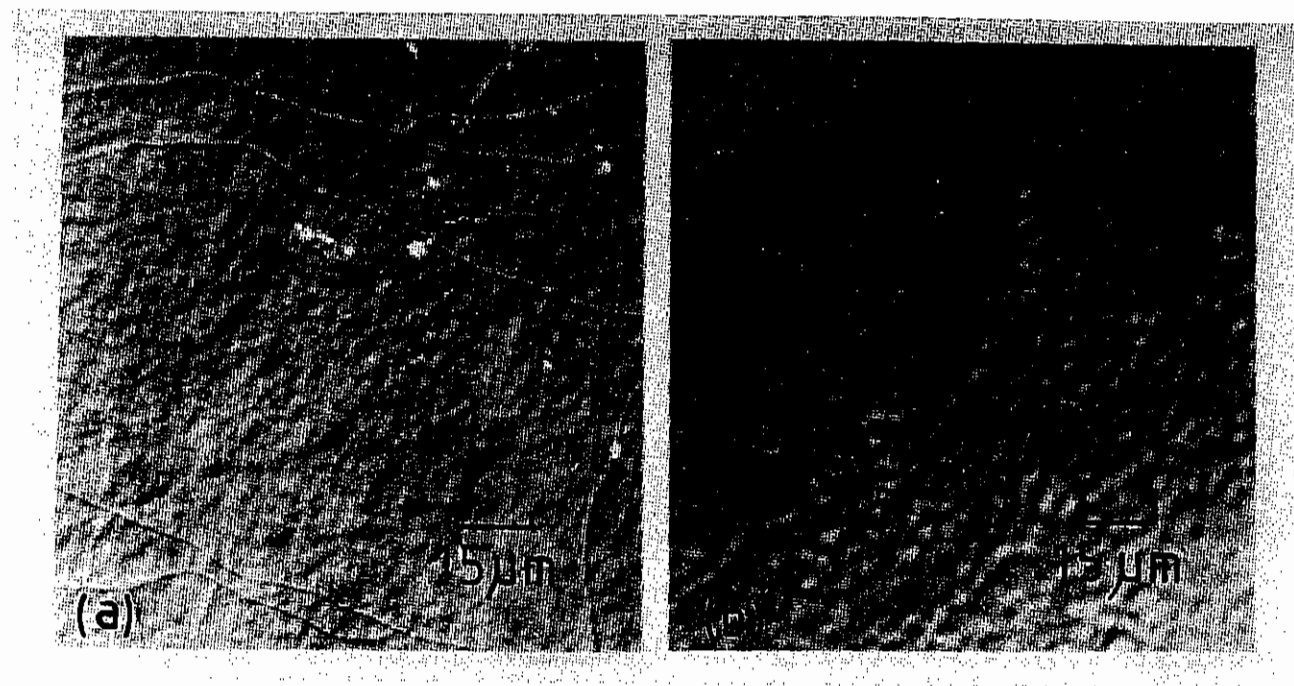


Fig. 3.14 Normarski micrographs of $\sim 0.2 \mu\text{m}$ thick sputtered films on GaAs annealed at 1.3 J cm^{-2} under 150 psig of Ar (a) XeCl radiation (b) XeF radiation.

densities. For the 1.0 μm films on Si, RHEED patterns similar to Figure 3.13(a) were observed up to $\sim 2.5 \text{ J cm}^{-2}$. However, below $\sim 1 \text{ J cm}^{-2}$ evidence of initial fibre texture in the films was also present.

Nomarski micrographs of the 0.2 and 1.0 μm thick films on Si annealed using XeF or XeCl radiation reveal varying amounts of surface rippling, often with a cellular structure. On the other hand, the 0.2 μm films on GaAs substrates exhibit severe cracking after XeCl anneals. This deleterious effect is however absent when XeF radiation is used. These latter observations are illustrated by the micrographs of Figure 3.14.

Finally, Figure 3.15 shows the results of SIMS profile measurements on Mn implanted films before and after laser and thermal annealing. The data was obtained on the specimens of Figure 3.11 and shows the flattening of the initially Gaussian Mn profile in the first $\sim 200 \text{ \AA}$ of the film after laser annealing. This supports the hypothesis of laser-induced surface melting and dopant re-distribution within this $\sim 2000 \text{ \AA}$ thick molten region.

3.5.3 Discussion

Our results show that XeCl and XeF pulsed laser annealing of ZnS films on Si or GaAs results in polycrystalline regrowth under the conditions described above. For the 1.0 μm sputtered specimens, regrowth probably occurs from the underlying polycrystalline film as it is unlikely that melt front penetration to the underlying substrate occurs (see Figure 3.15). Similarly the Mn-implanted single crystal films probably regrow from implantation damaged layers as Figure 3.15 suggests. Epitaxial layers implanted at lower energies are being investigated to clarify the situation. In contrast, for 0.2 μm thick specimens, it is likely that melt-front penetration to, and regrowth from, the underlying substrates takes place, particularly above $\sim 2 \text{ J cm}^{-2}$. Consequently these regrown layers exhibit a purely cubic structure characteristic of the substrates. The fact that both the ZnS/Si and ZnS/GaAs film/substrate combinations have either imperfect thermal or lattice matches (see Table 3.01) may well account for the observed polycrystalline nature of these cubic-phase

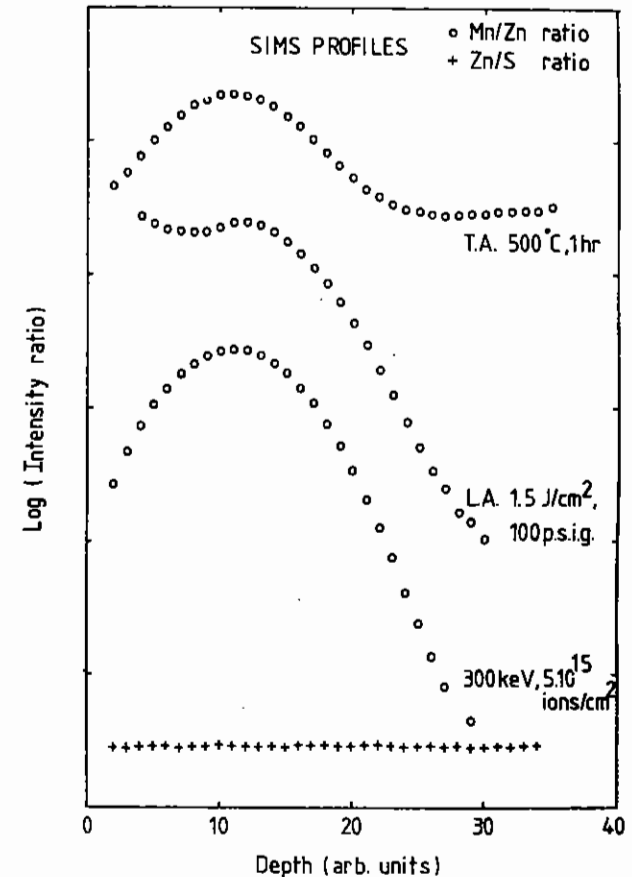


Fig. 3.15 SIMS profiles on Mn-implanted films of Figure 3.11 before and after laser and thermal annealing.

TABLE 3.01

Thermal expansion coefficients and lattice parameters for Si, GaAs and ZnS

Semi conductor	Lattice parameter (Å)	Expansion coefficient at 300 K ($10^{-6}K^{-1}$)
Si	5.431	2.33
GaS	5.654	5.8
ZnS cubic hexagonal	5.409 a = 3.814, c = 6.258	~ 6.2

TABLE 3.02

Summary of RHEED Analyses on sputtered films annealed under ~ 150 psig of Ar.

Sample Type	Substrate Cleaning	ZnS thickness μm		RHEED RESULTS	
				XeCl	XeF
			$\sim 0.6-1.8 J cm^{-2}$	$> 2 J cm^{-2}$	$\sim 1.3-1.8 J cm^{-2}$
ZnS:Mn/Si	Chemical only	~ 0.2	PR	PS	PR
ZnS/Si	Ar-ion etching	~ 0.2	PR	PS	PR
ZnS/Si	"	~ 1.0	PT	PR	PR
ZnS/GaAs	Chemical only	~ 0.2	PR	PS	PR

PR = Sharp Polycrystalline rings consistent with cubic + hexagonal phase material.

PS = Polycrystalline rings + spots; purely cubic phase material.

PT = As for PR but initial fibre texture of films also present at $< 1 J cm^{-2}$.

regrown layers; in any case it could contribute, with other possible factors, to the lack of epitaxial regrowth. It should also be noted that film crystallinity in these specimens (and also in the other cases studied) does not appear to depend on interface cleanliness or Mn doping of the films.

Of the other points to note is the observation of a possible hexagonal phase in the films after laser annealing, as described above. ZnS exhibits two stable allotropes, cubic and hexagonal. The latter is the stable form at high temperatures with the hexagonal to cubic transition being reported to occur at $\sim 1150^{\circ}C$. Thus competition between the two allotropes may possibly occur during regrowth after melting (the melting point of ZnS being $\sim 1830^{\circ}C$ under pressure). For regrowth primarily within the films (as for the thick specimens and for thinner samples annealed at relatively low powers) this may lead to loss of crystallinity and mixed cubic plus hexagonal material as observed. For regrowth from the underlying cubic substrates, however, the hexagonal phase may well be suppressed and purely cubic material obtained as discussed above. However further work is required to confirm this and the identification of the additional structure in the RHEED patterns with a hexagonal phase, rather than with twinning and double diffraction effects. This is underway.

The room temperature absorption coefficient of ZnS films at 351 nm and 308 nm has values of $\sim 4 \times 10^4 cm^{-1}$ and $\sim 1 \times 10^5 cm^{-1}$, respectively. Thus, one might expect the cooling rates after laser annealing to be greater for the XeCl anneals. However no significant effects due to this possibility have been observed in the crystal structure of the regrown material, although it could perhaps explain why severe film cracking occurs on GaAs substrates when using XeCl radiation but not when the XeF line is used (see Figure 3.14). Finally the cellular microstructure seen in many of the annealed films, e.g. Figure 3.12(a), may also be due to microcracking, at least in the thin films. However, other explanations are also possible and a fuller study using transmission electron microscopy (TEM) is underway.

In conclusion the programme of research will continue with further studies on film structure using TEM measurements. However, the effort will be concentrated on studying the electroluminescent properties of the annealed layers, together with their electrical properties using various dopant ions.

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C.B.Edwards (RAL).

3.6 Excimer laser lithography

Recently a series of preliminary experiments have been conducted at the Rutherford Appleton Laboratory UV Laser Facility. The work includes an investigation of excimer laser lithography at 193 nm. Using contact printing, gold zone plate masks and P.M.M.A resist, feature sizes of $< 0.5 \mu\text{m}$ with excellent wall profiles have been demonstrated using a single ~ 15 ns pulse of $\sim 80\text{mJ}/\text{cm}^2$ energy density.

3.6.1 Introduction

The resolution, w , obtained in optical contact or proximity lithography scales as $\sqrt{\lambda}$. For higher resolution it is therefore advantageous to use the shortest wavelength of light available. Traditionally deep - UV deuterium (D_2) or Xe-Hg arc lamps have provided the shortest wavelengths of 200 - 300nm. However, these sources have low output powers at these wavelengths ($< 50\text{mW}$). Hence for deep-UV resists such as polymethylmethacrylate (PMMA) which is sensitive for $\lambda < 0.25 \mu\text{m}$ long exposure times typically of several minutes are required. We have recently investigated the use of a pulsed rare gas halide ArF excimer which due to its high power output at $\lambda=193\text{nm}$ ($\sim 20 \text{ MW}$) can reduce the exposure time to ~ 15 nanoseconds. This remarkable decrease in exposure time essentially eliminates problems due to mechanical vibrations in the system as well as offering the potential of a much higher throughput of fabricated devices. In addition, since the photosensitivity of the resist is no longer of paramount importance, the use of an excimer laser source will allow the optimization of other resist properties such as higher image resolution

and wet-etch characteristics. Since excimer lasers radiate very little visible or infra-red radiation special precautions taken when using conventional sources to prevent distortion of the mask and wafer should not be necessary.

Thus there are several significant potential advantages of using excimer laser sources to replace conventional lamps. It is the purpose of the experiments reported in this work to investigate and demonstrate some of these advantages.

3.6.2 Experiments

An Argon Fluoride (ArF) 193 nm Lambda Physik EMG 201 excimer laser was used in these experiments to investigate the required energy density and resolution obtainable in PMMA resist for contact lithography.

The resist used was a $\sim 1\mu$ layer of cross-linked PMMA spun on a borosilicate glass (Pyrex) substrate coated with a chromium layer $850 \text{ \AA} + 50 \text{ \AA}$ thick. This provided an overall master-grade quality structure flat to $\sim 3 \mu\text{m}$ over a two inch diameter circle. Gold zone plate foils which were used as the masks were sandwiched between the PMMA resist and a quartz plate flat to $\sim 2 \mu\text{m}$ so that fairly hard contact was provided between mask and resist.

The mask/resist sandwich was mounted normal to the laser beam. A card with a suitable sized aperture was fixed to the quartz plate to single out one zone plate at a time per illumination.

To provide a range of energy densities at the resist the unfocused $50\text{mJ}/\text{cm}^2$ output of the laser was either increased by using a 50cm focal length quartz lens, or decreased by using a Rh6G + water absorption cell and varying the concentration of Rh6G. The lens or cell was placed directly in front of the laser. The energy density incident on the quartz plate was measured by a GenTec ED 500 Joule meter placed at a position of the plate, and the area of the beam was measured from burn patterns on UV photosensitive paper.

Each zone plate was exposed by a single laser shot, of $\sim 15\text{ns}$ duration, with a dosage varying from $\sim 1\text{mJ}/\text{cm}^2$ to $325\text{mJ}/\text{cm}^2$. The exposed plates were then developed using the 'spray' technique - 240s of developemnt is MIBK (methyl isobutyl ketone) followed by 20s rinsing in IPA (isopropyl alchol), and then 20s drying with N_2 gas. The developed plates were then examined using a CAMSCAN S.E.M.

3.6.3 Results

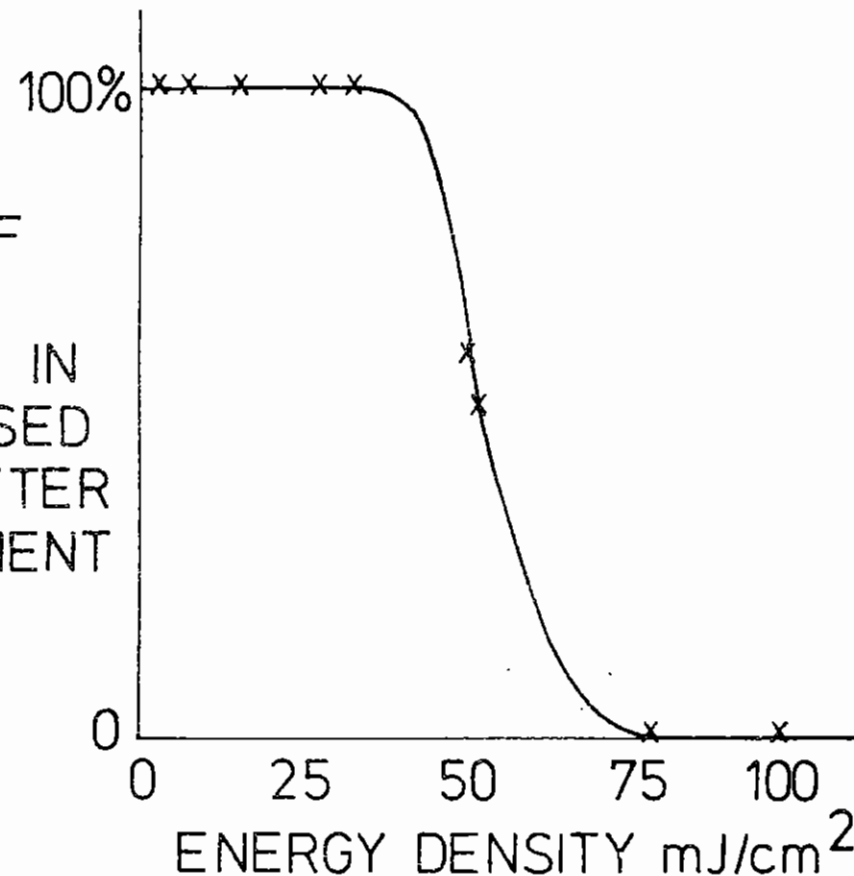
The results of these experiments are summarized in Fig. 3.16 which shows the amount of resist remaining in the exposed region after development, versus energy density. The threshold for PMMA resist removal during development is seen to be $\sim 30\text{mJ}/\text{cm}^2$, with an increasing proportion of the exposed resist being removed as the energy density increased up to $80\text{mJ}/\text{cm}^2$ and above which removal was complete.

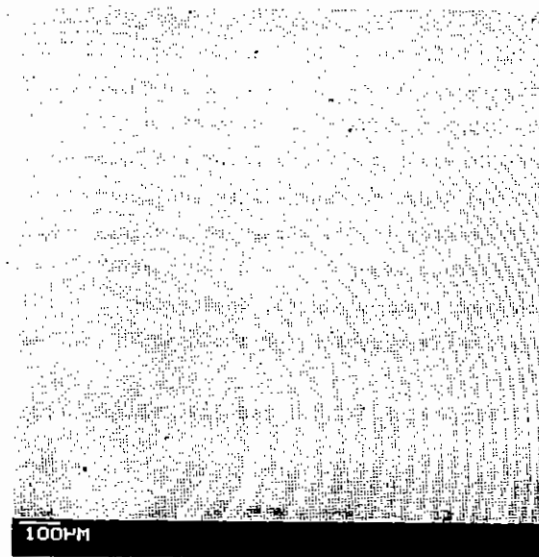
We found that the underdeveloped regions were patterned by diffraction/interference fringes within the aperture boundaries. Because of interference effects some areas received less energy density than others and so were less well developed. However, above $80\text{mJ}/\text{cm}^2$ there was sufficient dosage to saturate the resist over the whole aperture so that diffraction/interference effects were no longer observable.

The scanning electron micrographs shown in Fig. 3.17 are of the exposed photoresist using the Fresnel zone plate mask shown in Fig. 3.17(A) and using a fluence of $100\text{mJ}/\text{cm}^2$. The photographs shown in Figs. 3.17(B)-(D) were taken at an angle of 75° and show the relatively sharp wall profiles, and sub $0.5\text{ }\mu\text{m}$ lines.

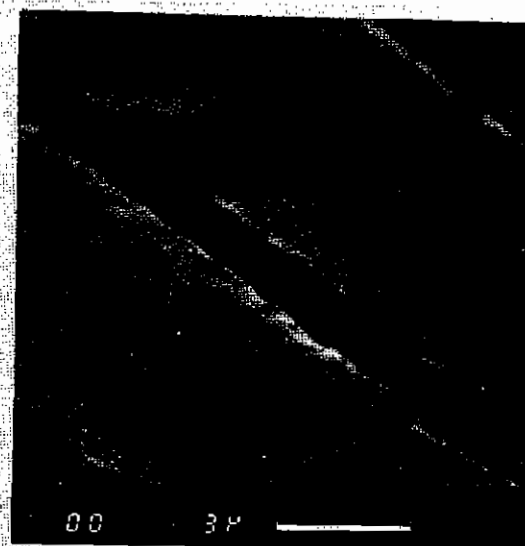
We conclude therefore from these preliminary experiments that with an ArF laser operating at 193nm it is possible to expose feature sizes of $< 0.5\text{ }\mu\text{m}$ with good image definition using a PMMA photoresist. Further experiments with shortly be carried out using finer line masks in order to determine the ultimate resolution of this type of contact printing.

FIG. 3.16
AMOUNT OF
RESIST
REMAINING IN
THE EXPOSED
REGION AFTER
DEVELOPMENT

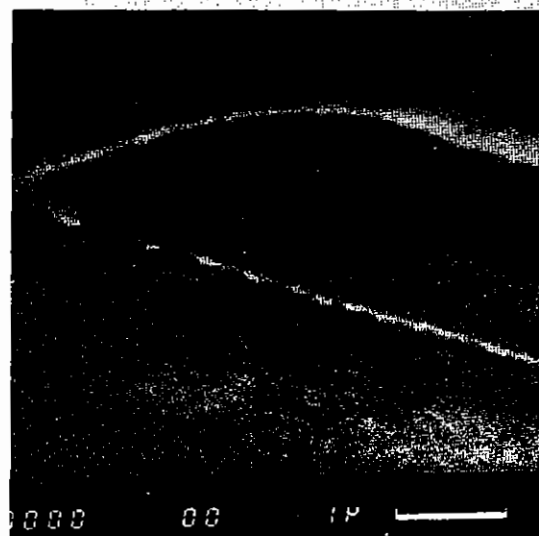




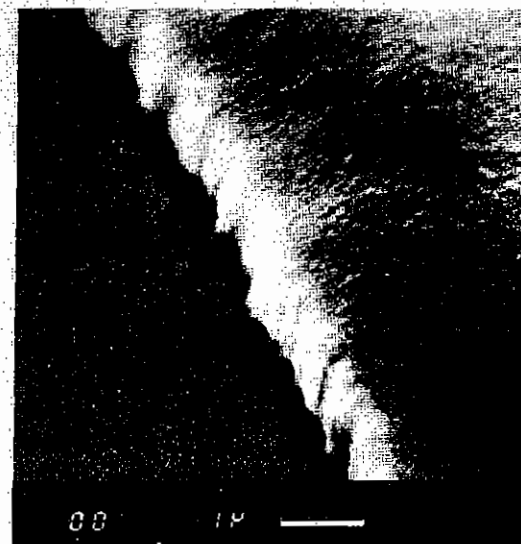
(A)



(B)



(C)



(D)

Fig. 3.17 Scanning electron micrographs of exposed photoresist using Fresnel zone plate mask shown in (A). Exposure is 100 mJ/cm^2 of ArF laser radiation.

3.7 UV Thomson Scattering on Non-Debye Discharge Plasmas and Associated Lineshape Studies

Preliminary experiments on UV Thomson scattering in the dense ($n_e = 10^{19} \text{ cm}^{-3}$), non-Debye ($N_p \sim 1$) z - pinch plasma device developed at Imperial College were carried out using the KrF 248 nm excimer laser system at the Rutherford UV Radiation Facility. Although the EMG 201 laser employed had not been line-narrowed at that time (16-22nd December 1982) initial experiments were carried out, including UV alignment of the 90° experimental scattering geometry and synchronisation of the z-pinch and laser pulses. Rayleigh scattering of the broadband 248nm radiation was also performed in atmospheres of argon and nitrogen, and the scattered signal was found to increase linearly with gas pressure (as expected). The extrapolated signal at zero pressure gave a measure of the stray light level, which when compared to the expected signal at $n_e = 10^{18} \text{ cm}^{-3}$ (obtained from the Thomson to Rayleigh cross-section ratio) was only 3%. Since in this high ($3 < \alpha < 6$) scattering experiment most of the scattered signal (>90%) will be in the central ion feature, this level of stray light is a negligibly small amount. The chief density diagnostic will be the separation of the electron feature from the ion feature (some tens of Å) and when the line narrowed excimer laser becomes available at the UVRF the visibility of the electron feature above the plasma background will be tested in the complete Thomson scattering experiment.

During the second scheduled period (7th to 20th February, 1983) experiments were performed on the production of tunable, coherent VUV radiation for use in a future absorption experiment on the z-pinch plasma at Lyman- α . The FL2001 excimer pumped dye laser was used to generate the fundamental radiation at $\sim 3650\text{\AA}$, and pulses of 4mJ in 4ns (giving a conversion efficiency from the excimer laser of 2.8%) were produced. The fundamental beam at 3650Å was focused into a krypton tripling cell developed at I.C. and generation of VUV Lyman α radiation was achieved through 3-photon upconversion of the fundamental. The Lyman - α radiation

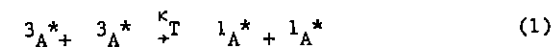
generated was at least an order of magnitude greater than had previously been produced at I.C. using a much smaller dye laser system, and varied within a factor of 3 except near the limits of the maximum theoretical range possible ($\pm 15 \text{ \AA}$ from Lyman- α obtained by altering the pressure of the krypton phase matching medium). A small newly constructed monochromator used to monitor the Lyman- α radiation was successfully tested and was shown to discriminate adequately between Lyman- α radiation and the much greater ($\sim 10^6$) fundamental signal. Measurement of the plasma emission at Lyman- α showed the intensity to be of the same order as the coherent Lyman- α signal, indicating that the Lyman- α absorption experiments proposed for the next scheduled period should be feasible.

D.D.Burgess, K.Baldwin and J.P.Marangos (Imperial College)

3.8 Trial Experiments

3.8.1 UV Saturable Absorbers

The recent rapid developments in the area of high power excimer lasers has lead to an urgent requirement for ultra-violet saturable absorbers suitable for pulse shape control and optical isolation in oscillator-amplifier chains. For KrF based excimer system operating at 248 nm the aromatic molecule anthracene is a promising candidate due to its extremely high absorption cross section at 248 nm. Anthracene solutions have the added advantage that for a fixed low power transmission level the saturation characteristics and absorption recovery time can be varied to some extent either by addition of excited state quenchers (e.g. oxygen) or by varying the anthracene concentration-absorption path length combination used leading to a change in the rate of triplet-triplet annihilation as shown in reaction (1).



The saturation characteristics of a series of anthracene-hexane solutions were investigated using the $\sim 1 \text{ J}$ pulse of 20 ns duration from a Lambda

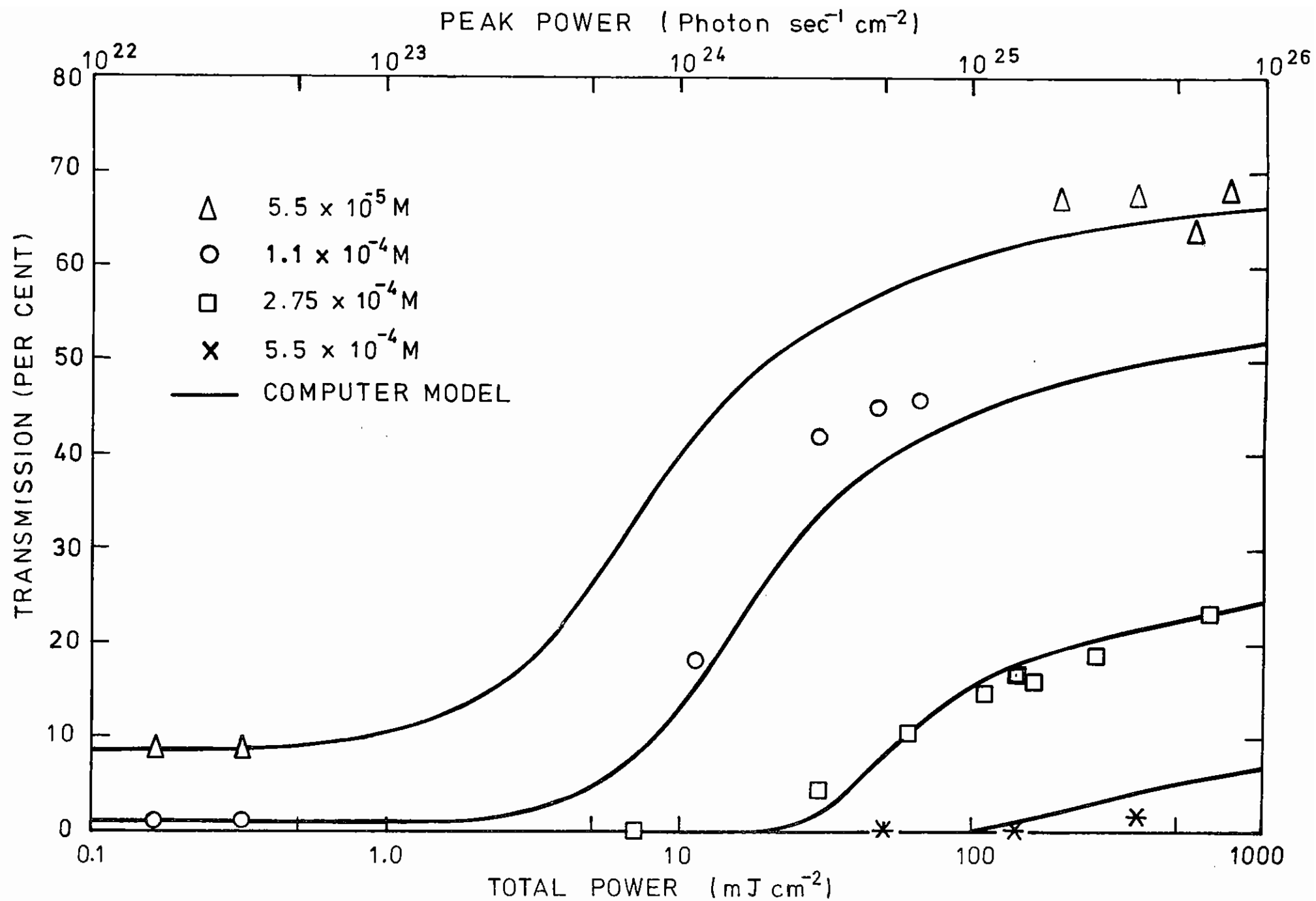


Fig. 3.18 Saturation characteristics of anthracene - hexane solutions for KrF laser radiation.

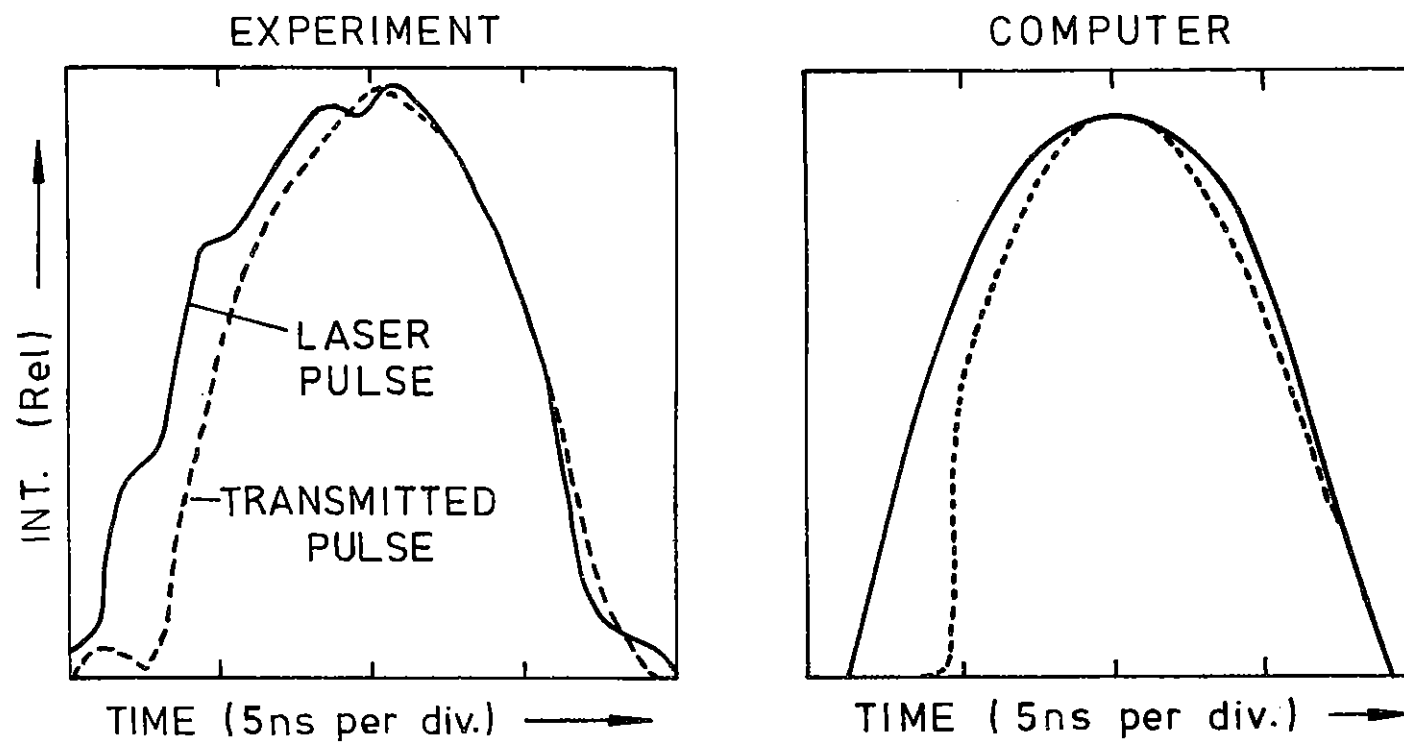
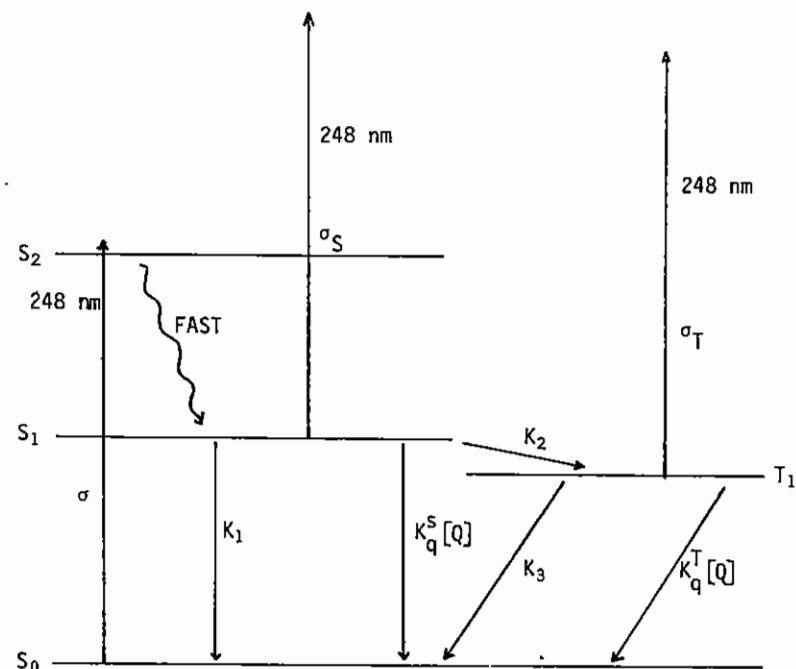


Fig. 3.19 Temporal profiles of incident and transmitted laser pulses for anthracene ~ hexane saturation experiment.



σ	Ground state absorption cross section at 248 nm	$7 \times 10^{-16} \text{ cm}^2$
σ_S	Excited singlet absorption cross section at 248 nm	$7 \times 10^{-17} \text{ cm}^2*$
σ_T	Excited triplet absorption cross section at 248 nm	$7 \times 10^{-17} \text{ cm}^2*$
K_1	Rate constant for radiative and non radiative decay of S_1 to S_0	$6 \times 10^7 \text{ sec}^{-1}$
K_2	Rate constant for intersystem crossing	$1.4 \times 10^8 \text{ sec}^{-1}$
K_3	Rate constant for non-radiative decay of T_1 to S_0	$< 10^4 \text{ sec}^{-1}$
K_q^S	Rate constant for oxygen quenching of S_1	$2.5 \times 10^{10} \text{ sec}^{-1}$
K_q^T	Rate constant for oxygen quenching of T_1	$0.29 \times 10^{10} \text{ sec}^{-1}$
K_T	Rate constant for reaction (1)	$7 \times 10^9 \text{ M}^{-1} \text{ sec}^{-1}$
$[Q]$	Oxygen concentration in aerated hexane	$2.3 \times 10^{-3} \text{ M}$

*Values which give best fit to experimental results.

Physik KrF excimer laser. Since the power output of the laser could not be varied directly without changing the temporal profile the output beam was focused with a quartz lens and the power incident on the 1 mm path length saturable absorber cell was varied by changing the position of the cell relative to the beam focus. The results, shown in Figure 3.18, were obtained by measuring the pulse energy transmitted using a calibrated calorimeter. Temporal profiles of the incident and transmitted pulses were monitored on UV photodiodes and typical profiles for a 140 mJ cm^{-2} pulse incident on a $2.75 \times 10^{-4} \text{ M}$ solution are shown in Figure 3.19. In figure 3.19 the transmitted pulse profiles have been normalised to the same peak value as that of the incident pulse.

The computer simulations of the saturation process, shown in figures 3.18 and 3.19, were based on the simple energy level scheme shown in figure 3.20 and available literature values for the appropriate rate constants. Unfortunately accurate values of σ_S and σ_T at 248 nm are not available so these parameters were varied to obtain the best fit to the experimental results. Experiments are at present underway to determine accurate value for σ_S and σ_T at 248 nm and also to investigate the influence of multiphoton absorption processes at higher power densities (i.e. $> 10^{26} \text{ photons sec}^{-1} \text{ cm}^{-2}$).

T.Morrow and D.Denvir (Queen's University, Belfast)

3.8.2 Laser induced ionisation in proportional counters

The problem of laser induced ionisation in proportional counters has received much attention during the last two or three years. The impetus for this study has come from two disparate fields of research. Firstly the need to calibrate multiwire drift chambers can be satisfied by simulating particle tracks with laser induced ionisation(3.21) –(3.26). Secondly the ultrasensitive analytic technique of multiphoton resonant ionisation often requires the use of proportional or ionisation chambers (3.27) – (3.29). Much of the work carried out to date, has used nitrogen lasers operating at 337 nm with subnanosecond pulses.

Fig. 3.20 Anthracene energy level scheme used for computer simulation of saturation process.

It was felt that a systematic study of the ionisation caused by different types of laser, operating at different wavelengths and pulse lengths was necessary to establish the best kind of laser for this work. The present study has compared the ionisation induced by a nitrogen laser at 337 nm and pulse length 0.3 ns with:-

- a) An excimer laser using XeF (351nm), XeCl (308), KrF (248) and ArF (193 nm) with pulse length about 15 ns.
- b) An excimer pumped dye system operating at around 500 nm with pulse length about 5 ns.
- (a) & b) were carried out using the RAL UV Radiation Facility.]
- c) A Nd- glass laser operating at its fundamental frequency (1060 nm), doubled (530), tripled (354) and quadrupled (265) with pulse lengths about 25 ns.

A proportional counter, carefully end corrected, shown in Figure 3.21, was used and ionisation was induced at all wavelengths. Fig 3.22 shows the ionisation induced with the XeF excimer at 351 nm. In the ArF/10% CH₄ filling of the counter. It was found that the specific intensity of the ionisation increased considerably when the laser wavelength was decreased from 350 nm to 193 nm. It is hoped to carry out further investigations using frequency doubling crystals between 200 - 300 nm and also varying the pulse lengths of the lasers.

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3.8.3 Fluorescence of Liquid Phase Toluene Excited By a KrF Laser

The unique combination of high pulse power and a variety of output wavelengths ranging from XeF (351 nm) to ArF (193 nm) available from excimer lasers makes such systems invaluable in the studies of:-

- a) two photon absorption processes
- b) examination of basic photophysical properties of solvents as a function of wavelength.

We have used the KrF excimer laser at the RAL UV Radiation Facility for an

initial part of a study of the absorption and emission properties of toluene as a function of laser power and wavelength. It is well known that aromatic solvents can be excited in regions above their absorption bands by two photon processes, though the relative efficiency of these processes compared with monophotonic absorption is not known.

De-aerated samples of toluene were excited by the 249 nm line of a KrF laser (pulse width = 20 ns) in the experimental arrangement shown in Fig. 3.23. Variation in the laser intensity was achieved by using dilute solutions of toluene in hexane as neutral density filters for the laser beam. The laser intensity was monitored via the output of a calibrated GENTEC ED 200 calorimeter using a reflected percentage of the beam. The fluorescence output was monitored using an OMA 2 vidicon. The total fluorescence yield was obtained by integration of the whole fluorescence spectrum (after subtraction of a background). The total counts per channel was kept to well below maximum values and two detector voltages were used in an attempt to maintain linearity (specified by EG and G to be $\pm 2\%$).

The data obtained is given in Fig. 3.24. The fluorescence yield for both detector voltages may be reasonably fitted to an equation of the type $Y = AI + BI^2$ where I is the laser intensity and A and B are constants. The solid lines are for best fits to the data, A and B differing only in a factor of 0.39 to account for the different detector sensitivities.

$$8.0 \text{ kV } A = 0.65 \quad B = 0.0125$$

$$10.0 \text{ kV } A = 1.67 \quad B = 0.0346$$

The non linear production of toluene fluorescence may be due to one of two processes.

- 1) singlet states produced by triplet-triplet annihilation
- 2) two photon excitation to higher excited states which also lead to fluorescence via internal conversion rather than ionisation.

Since the time profiles of these two possibilities will differ, it is hoped to distinguish them by further time resolved measurements. Use of

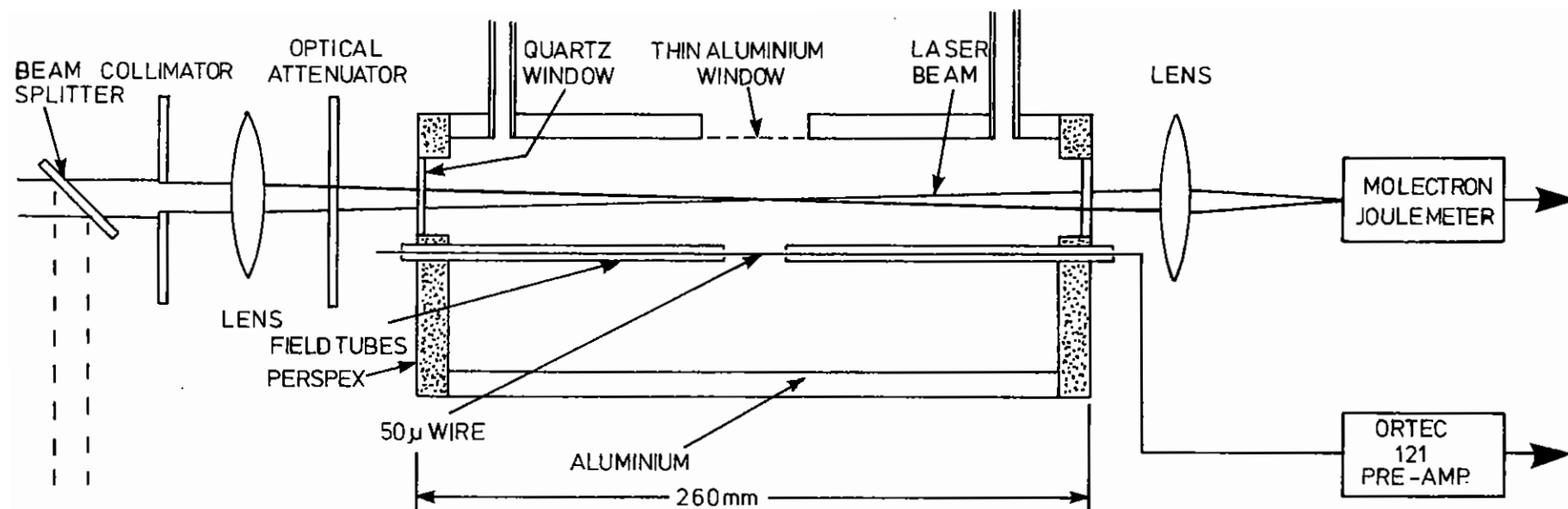


Fig. 3.21 Proportional counter.

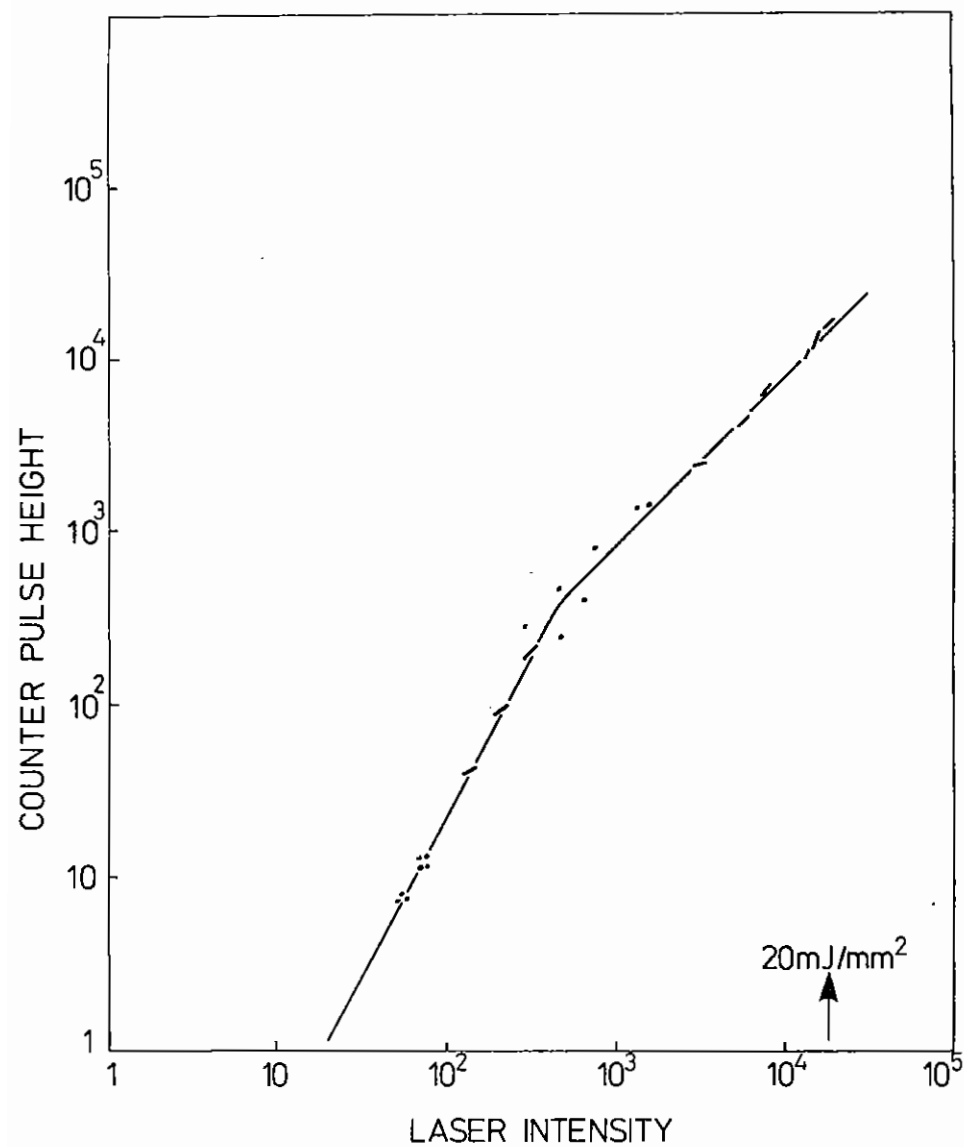


Fig. 3.22 Ionisation induced by XeF at 351 nm. Counter filled with 10% CH_4 in Ar.

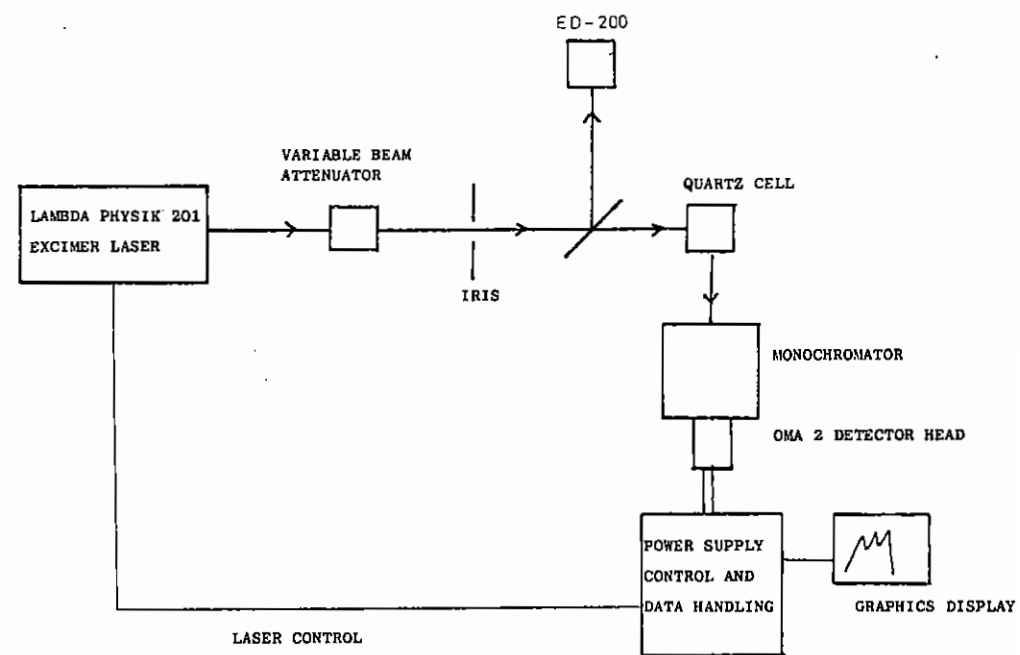


Fig. 3.23 Experimental arrangement.

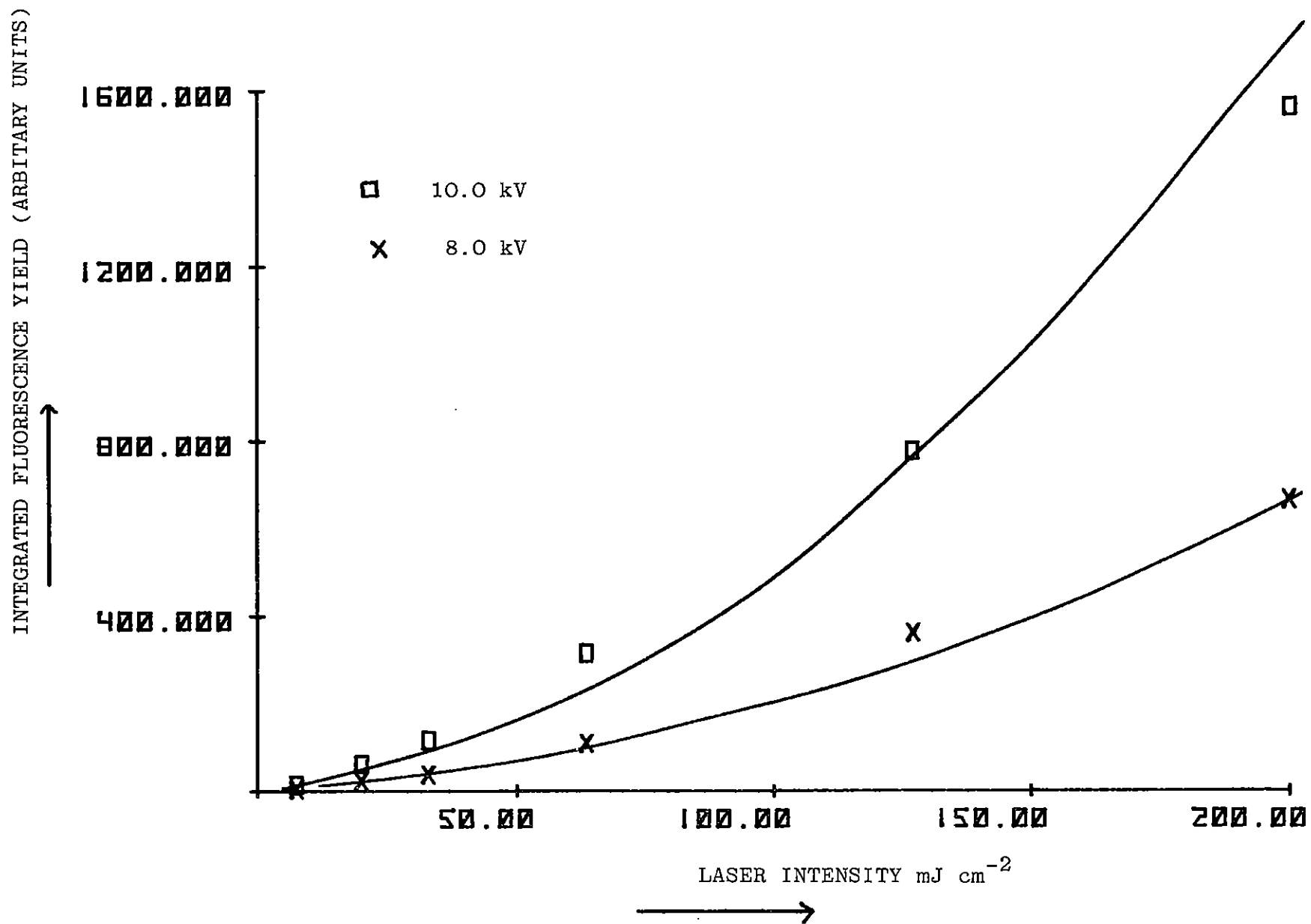


Fig. 3.24 Fluorescence yields of toluene as a function of incident KrF laser energy density for two different detector voltages.

other laser wavelengths will enable the efficiencies of the processes to be compared.

R.B.Cundall and J.T.Richards (Department of Biochemistry, University of Salford).

Chapter 3

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INDEX

- 4.1 INTRODUCTION page 4.1
- 4.2 SIMULTANEOUS OBSERVATIONS OF $\omega_0/2$ AND $3\omega_0/2$ HARMONIC EMISSION
FROM 526 nm LASER PRODUCED PLASMAS page 4.1
- 4.3 INSTABILITY OF FILAMENTATION AND BACKSCATTER IN LASER PRODUCED
PLASMAS page 4.5
- 4.4 FILAMENTATION INSTABILITY IN LASER PRODUCED PLASMAS page 4.9

REFERENCES page 4.15

CHAPTER EDITOR: R G Evans

4.1 Introduction

This year has seen two notable new results together with the extension of the study of plasma "jets" that has been reported in previous years. The simultaneous observation of the spectra of $\omega_0/2$ and $3\omega_0/2$ radiation has shown that the mechanism for the generation of $3\omega_0/2$ light is the photon plus plasmon process rather than the three plasmon process. The same experiment suggests that the $\omega_0/2$ emission is produced by the analogous process of Raman down conversion but this process can only phase match in an inhomogeneous plasma due to refraction of the electron plasma waves.

The observations of the spatial structure of the "burn through" image in thin foil targets shows the first evidence for unstable light filaments, theoretical mechanisms for this have been suggested for many years and some new ideas are also presented in this chapter.

Thin foil targets are expected to be a continuing feature of the laser plasma interaction programme since they offer a more controllable evolution of the density profile than is available with solid targets.

4.2 Simultaneous Observations of $\omega_0/2$ and $3\omega_0/2$ Harmonic Emission from 526 nm Laser Produced Plasmas

Near the quarter critical density surface, two mechanisms exist which generate plasmons with frequency near $\omega_0/2$: the stimulated Raman scattering (SRS) instability (4.01) where the incident photon decays to give a scattered photon and a plasmon, and the two plasmon decay instability (4.02), where the incident photon decays to give two daughter plasmons both with frequency $\omega_0/2$. Scattered light with frequency near $\omega_0/2$ is produced directly by the SRS instability but may also be generated from plasmons created by the two-plasmon decay process, either by Stokes Raman scattering of the incident pump photon off these plasmons, or by

linear mode conversion (4.03). Similarly, scattered light with frequency near $3\omega_0/2$ can arise either as a result of the interaction of a photon with a plasmon, or the coalescence of three plasmons. Besides the intrinsic interest of these parametric instabilities they are also important as sources of very energetic electrons which can produce damaging levels of preheat in compression experiments.

Recent observations (4.04, 4.05) of $\omega_0/2$ harmonic emission spectra have shown that it possesses a double peaked structure similar to that previously reported in $3\omega_0/2$ harmonic spectra (4.06, 4.07). Hence, it was suggested (4.04, 4.05) that the generation mechanism for the $\omega_0/2$ harmonic light is connected with processes involving the two-plasmon instability. Scattered light at $\omega_0/2$ due to a two-plasmon instability re-emission process has been separately identified (4.08) in UV-laser produced plasmas. However, the mechanism by which the $\omega_0/2$ plasmons are converted to $\omega_0/2$ photons is still unclear. The former of the two processes suggested earlier would explain the similarity between the $\omega_0/2$ and $3\omega_0/2$ harmonic spectra, since they then represent the Stokes and anti-Stokes lines, respectively, of Raman scattering of photons of two-plasmon instability-generated plasmons, where the separation between the peaks for both harmonics is given by the frequency spectrum of the $\omega_0/2$ plasma waves. Hence conservation of energy requires that

$$\Delta\omega_{(1/2)} = \Delta\omega_{(3/2)}$$

$$\text{giving } \Delta\lambda_{(1/2)} = 9 \Delta\lambda_{(3/2)}$$

We present here new simultaneous observations of the backscattered $\omega_0/2$ and $3\omega_0/2$ harmonic emissions from nylon tube targets irradiated by 526 nm laser light to investigate the relationship between these two harmonic spectra.

The experimental layout is shown in Figure 4.01. The 1.05 μm laser light was frequency doubled by a KDP crystal and focused onto the target by an f/1 aspheric doublet lens. The residual fundamental light was attenuated by an HR 1 μm mirror in the incident beam path. Incident energies of

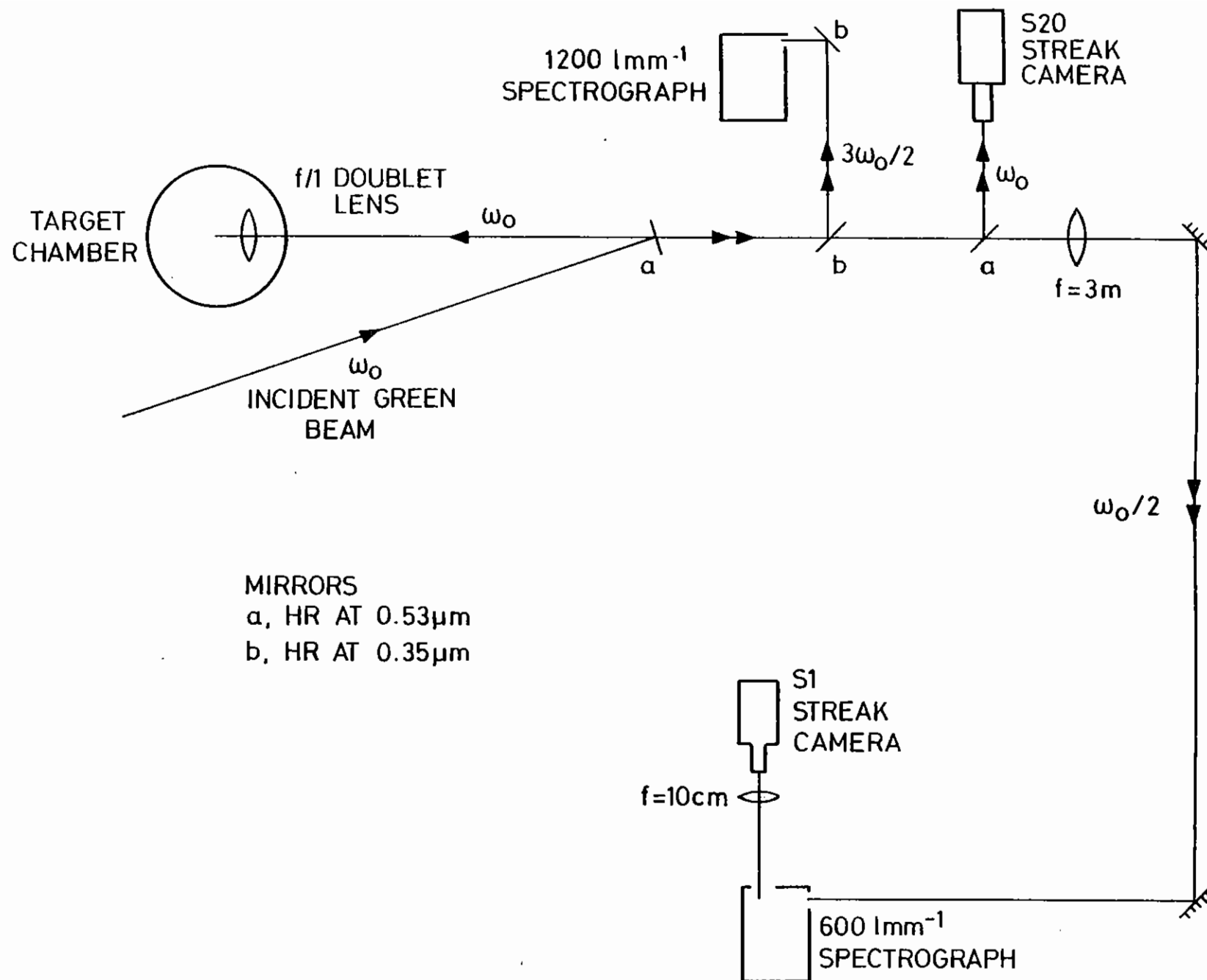


Fig.4.01 Experimental layout for simultaneous spectroscopy at $\omega_0/2$ and $3\omega_0/2$.

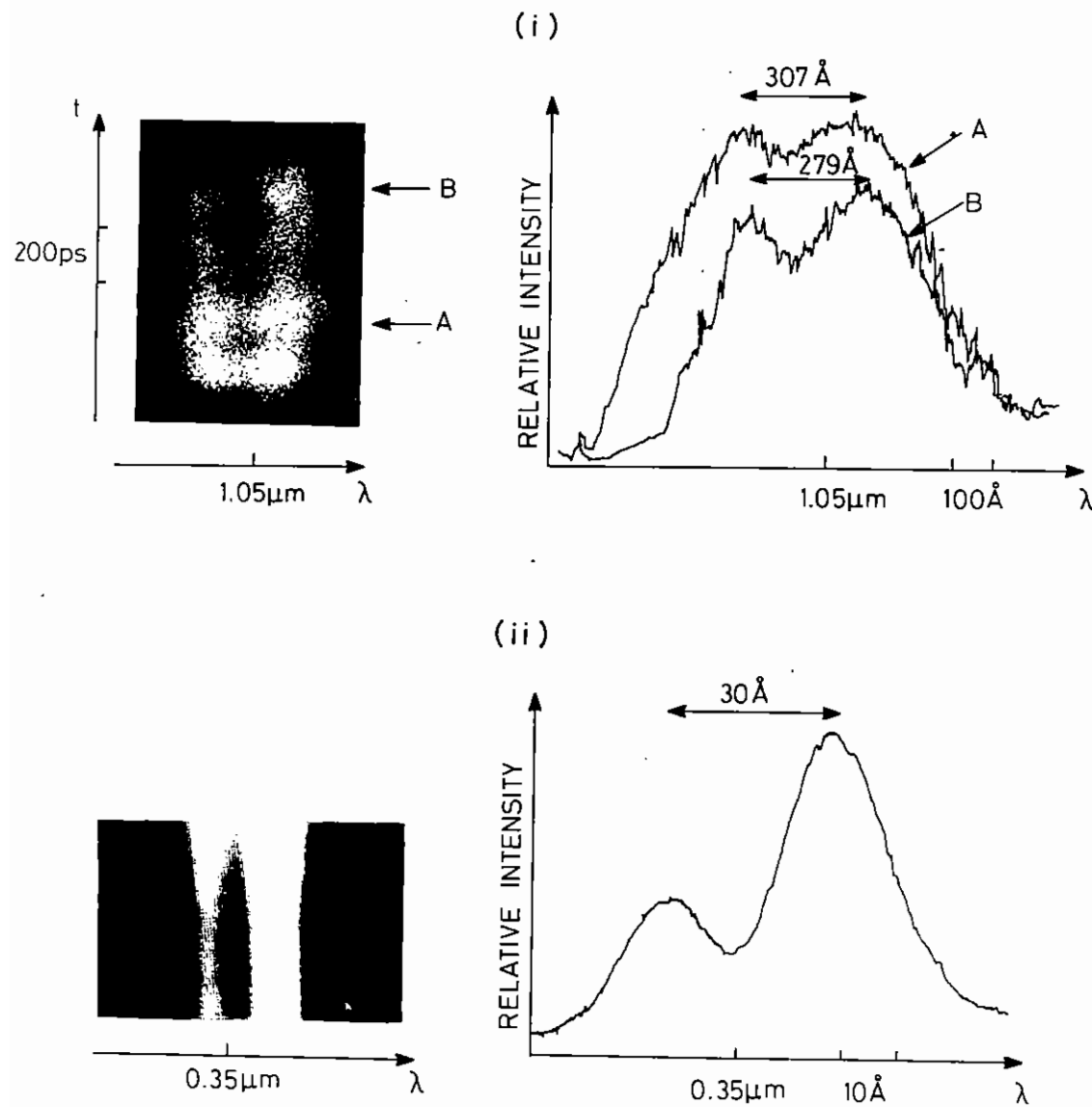


Fig 4.02 Simultaneous spectra at $\omega_0/2$ (time resolved) and $3 \omega_0/2$ (time integrated).

$\sim 16 \text{ J}$ on target were obtained for pulse lengths $\sim 800 \text{ ps}$.

Time resolved spectra of the $\omega_0/2$ harmonic light backscattered through the focusing lens were obtained by imaging the emission onto the slit of a grating spectrograph ($f = 0.5 \text{ m}$, $600 \text{ lines mm}^{-1}$ grating), optically coupled by an $f/1.4$ lens to an S-1 streak camera. High temporal resolution was not essential for this experiment and the streak camera was chiefly used for its high gain in this spectral region (due to its image intensifier). Temporal and spectral resolutions were $\sim 160 \text{ ps}$ and $\sim 50 \text{ \AA}$ respectively. Two red glass filters (transmitting for $\lambda > 7000 \text{ \AA}$) and an IR transmitting black glass filter were placed in front of the streak camera slit. The input optics were arranged to give a spectral window of $\sim 470 \text{ \AA}$ at the streak camera slit, centred at the nominal $\omega_0/2$ position. Recordings of the time resolved spectra were made on Ilford HP5 film. The spectral dispersion along the streak camera slit was calibrated by recording low energy $1.05 \mu\text{m}$ laser shots with the spectrograph at two different wavelength counter settings.

Time integrated spectra of the backscattered $3\omega_0/2$ harmonic light were resolved on an $f = 0.5 \text{ m}$, $1200 \text{ lines mm}^{-1}$ grating spectrograph and recorded on Ilford HP5 film. A UV-transmitting black glass filter was used in front of the spectrograph slit. The spectrograph was used in first, second and third orders to obtain suitable recording levels of light over the range investigated. Corresponding spectral resolutions were $\sim 16 \text{ \AA}$, $\sim 8 \text{ \AA}$ and $4 \text{ \AA}$. A mercury lamp was used to provide calibration lines for each shot.

A range of irradiances on target between 10^{14} and $2 \times 10^{15} \text{ W cm}^{-2}$ was obtained by varying the focal spot size between $30\text{--}100 \mu\text{m}$ diameter. Figure 4.02 shows the simultaneously observed $\omega_0/2$ and $3\omega_0/2$ harmonic spectra for a typical shot. The pulsed nature of the $\omega_0/2$ harmonic emission is due to mode beating in the main laser pulse. The $3\omega_0/2$ spectrum is time-integrated and therefore represents an "average" peak separation during the laser pulse, whereas the $\omega_0/2$ spectrum, despite the poor temporal resolution, shows a variation in peak separation with time. Since the peak separation is proportional to the electron

temperature (4.06, 4.07) the $\omega_0/2$ spectrum reflects the temporal variation in the temperature. Over the range investigated the separation between the blue-shifted and red-shifted peaks for the $\omega_0/2$ harmonic spectra was observed to be $\Delta\lambda_{(1/2)} \approx (206-330 \text{ \AA})$ and for the $3\omega_0/2$ harmonic emission to be $\Delta\lambda_{(3/2)} \approx (24-31 \text{ \AA})$.

These measurements show good agreement with the relationship $\Delta\lambda_{(1/2)} = 9\Delta\lambda_{(3/2)}$ predicted by a Stokes and anti-Stokes Raman scattering generation mechanism. However, in the absence of any other process, the observed light should be split symmetrically (4.07) about $\omega_0/2$ or $3\omega_0/2$ and the experimental observations show the $\omega_0/2$ mean is slightly blue-shifted, while the $3\omega_0/2$ is slightly red-shifted. Although this is consistent with the idea of Raman generation it suggests that either the plasmons from the two-plasmon instability have lost some energy, (perhaps due to ion acoustic turbulence or some other parametric decay instability) or alternatively, the asymmetric splitting of the spectra may be explained by k-matching considerations. Since the plasma waves have to mix with the pump waves to give scattered light at $\omega_0/2$ and $3\omega_0/2$, the plasma wave vectors $k_{p1,2}$ must satisfy both the scattered wave dispersion relation and the usual conservation of momentum and energy matching conditions, $k_s = k_p + k_o$ and $\omega_s = \omega_p + \omega_o$ (the subscripts s, p and o denote the scattered, plasma and pump wave respectively). The matching conditions cannot be satisfied simultaneously for both k_{p1} and k_{p2} (4.09), hence different pairs of two-plasmon decay waves contribute to each peak.

The observations do not support the generation of $3\omega_0/2$ harmonic emission by the coalescence of three plasmons. Neither was any evidence found for the production of $\omega_0/2$ scattered light and $\omega_0/2$ electron plasma waves by the SRS instability at the quarter critical density surface, although scattered light due to a convective SRS instability has been identified in plasma densities of $0.05 n_c$ (4.08) and $0.1 n_c$ (4.05).

In conclusion, the experimental observations have shown that

$\Delta\omega_{(1/2)} = \Delta\omega_{(3/2)}$, as predicted by the Stokes and anti-Stokes Raman generation mechanism. However, the shift in the mean wavelengths of the $\omega_0/2$ and $3\omega_0/2$ spectra have yet to be fully explained.

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4.3 Instability of Filamentation and Backscatter in Laser Produced Plasmas

The closely related instabilities of laser beam filamentation and stimulated Brillouin scattering (SBS) play a vital role in laser plasma interaction experiments. Filamentation may cause laser energy to be deposited locally causing non-uniformities in the ablation pressure and the production of hot electrons and hard X-rays. On the other hand SBS may cause a large fraction of the incident laser energy to be scattered back before it is absorbed. Both instabilities grow faster in large underdense plasmas and consequently become important in ablative implosion experiments where long laser pulses and large focal spots are used. Laser beam filamentation has recently been seen in laser produced plasmas by optical probing and harmonic emission observations (4.10-4.12). The self-focused filaments have been predicted to be unstable by several models (4.13-4.15) but up to now no experimental observations have been made that support this. We have performed two different kinds of experimental observations with very high spatial and temporal resolutions. Firstly observations taken on thin foil targets which become underdense during the laser pulse show that the laser light burning through is spatially and temporally modulated indicating that the filamentation process may be unstable. Secondly studies of time resolved back-scattered ω_0 spectra show modulations on a time scale of ~ 10 ps, very much faster than in any previous observations (4.16), and only a few ion wave periods in duration. This back-scattered laser light is turned off shortly after the plasma becomes underdense indicating that it originates from plasma densities close to critical.

Targets consisting of parylene foils of thickness $0.5\text{ }\mu\text{m}$ or wires of various atomic number materials (CH_2 , Al, W) were irradiated with a single laser beam ($\lambda = 0.53\text{ }\mu\text{m}$) of the VULCAN Nd glass laser system (4.17). The beam was focused onto target with an f/1 aspheric doublet lens. Beam energies of up to 30 Joules were used with a pulse duration of 0.9 ns (FWHM). Target irradiances between 2×10^{12} and $5 \times 10^{14}\text{ W cm}^{-2}$ were achieved. Figure 4.03 shows the experimental arrangement with which two different kinds of measurements were made. Firstly the radiation back-scattered through the main focusing lens was imaged onto the slit of a $0.5\text{ }\mu\text{m}$ focal length, 600 lines/mm grating spectrograph with a 20 x magnification. The output of the spectrograph was relayed with a 2 x magnification onto the input slit of an IMACON 675 streak camera. To improve the temporal resolution of the combined spectrograph-streak camera system a slit was inserted in front of the grating giving temporal and spectral resolutions of 5 ps and 2.5 Å respectively. The spatial resolution was $5\text{ }\mu\text{m}$. On some shots a biprism element was inserted in front of the streak camera slit which made it possible to obtain two spectra on the same image originating from different parts of the plasma (4.18). Simultaneously the temporal profile of the incident laser pulse was monitored by focusing a part of the incident laser beam onto the slit of the streak camera. Secondly when the target consisted of thin foils the rear surface of the target was imaged onto the slit of the streak camera in order to obtain a spatially and temporally resolved image of the laser beam burning through the target. Again the spatial resolution was $5\text{ }\mu\text{m}$.

Figure 4.04 shows a streak camera image of the spatial distribution of the light transmitted through a parylene foil $0.5\text{ }\mu\text{m}$ in thickness. The thickness of the foil was chosen such that it became underdense close to the peak of the laser pulse. The target was irradiated with an average intensity of $1.3 \times 10^{14}\text{ W cm}^{-2}$ in a focal spot of $50\text{ }\mu\text{m}$ diameter. As can be seen in Figure 4.04 the laser beam burns through in numerous small spots. Spots as small as $5\text{ }\mu\text{m}$ diameter can be seen which is the spatial resolution of the experimental arrangement. The envelope of the spots spreads in radius as the size of the hole through the foil increases with time. Each spot persists only for a few picoseconds. It is

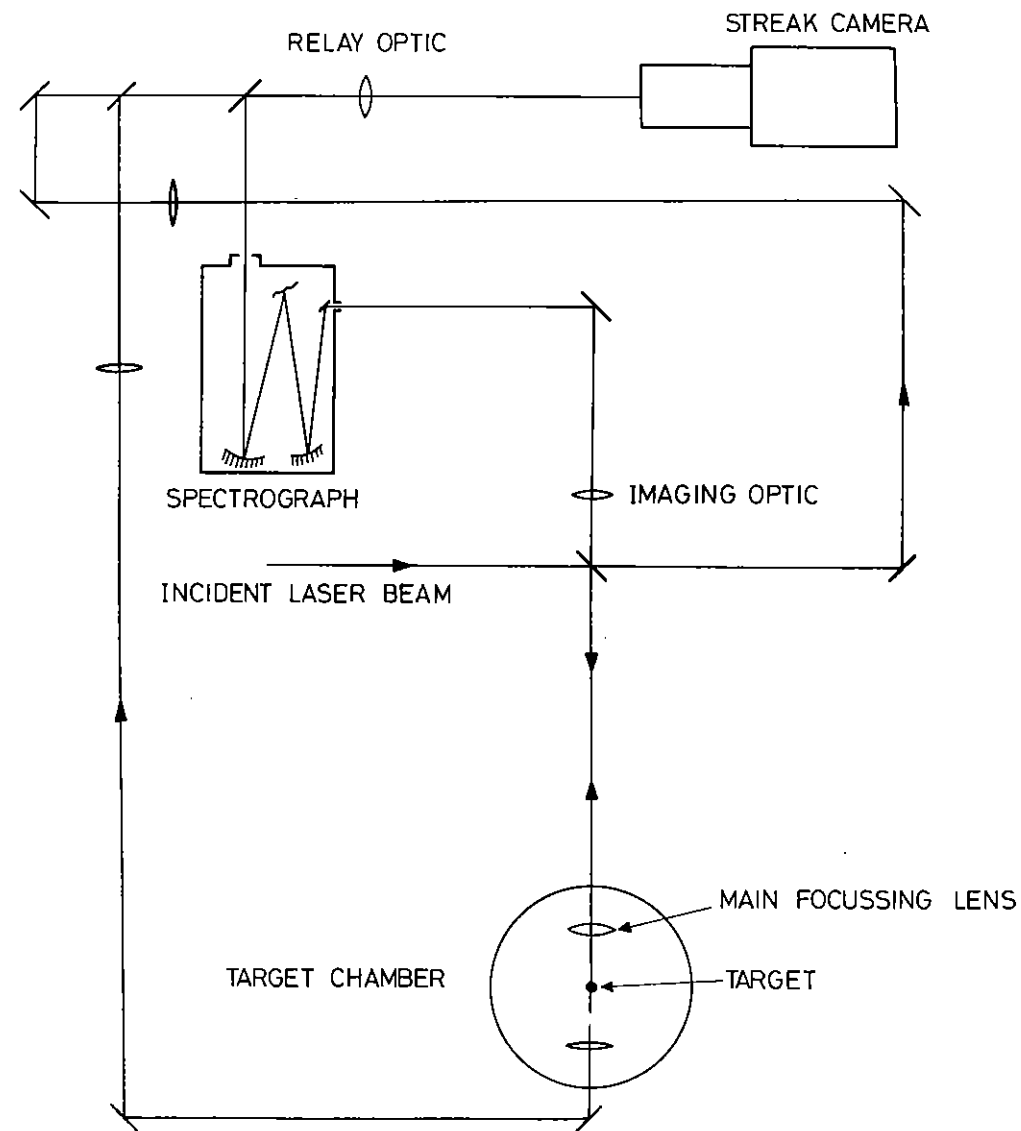
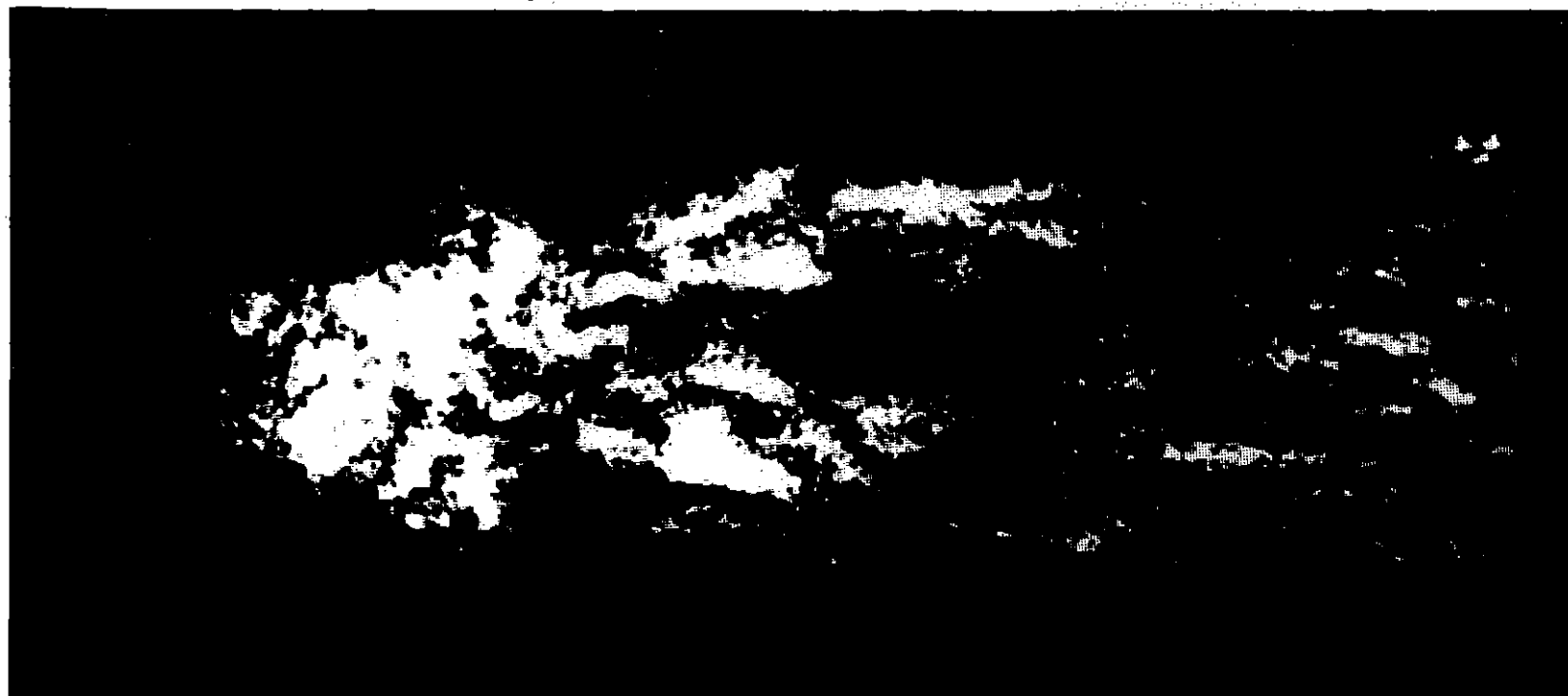


Fig. 4.03 Experimental layout for time resolved burnthrough imaging and backscatter spectroscopy.

space

50 μm



100 ps

time

Fig. 4.04. Time resolved beam through image on 0.5 μm parylene foil.

noticeable that at later times when the laser intensity and the plasma density are falling the transmitted spots are more stable in time. It is interesting to note that the spots are smaller in size than the hot spots in the incident laser beam measured in an equivalent target plane which were between 10 and 12 μm in diameter. This may be an indication of self-focusing effects within each filament. The temporal behaviour of the laser beam profiles in the target plane was also recorded with the streak camera. These observations clearly showed that the spatial beam distribution in the target plane did not change during the laser pulse.

Further observations of very rapid changes in the plasma were provided by time resolved spectra of the light back scattered into the focusing lens. Figure 4.05 shows a spectrum taken on an aluminium wire target together with the temporal profile of the incident laser pulse. The irradiance on target was $3.8 \times 10^{14} \text{ W cm}^{-2}$ and the laser beam was incident at 20° to the target normal. The back scattered spectrum shows very rapid and apparently random fluctuations with a time scale of 5 to 20 ps whereas no such modulations are observed in the incident laser pulse. Spectra taken on various materials such as plastic, aluminium or tungsten and on targets at various angles (0° to 45°) to the axis of the laser beam showed similar temporal behaviour. In addition spectra from two different parts of the plasma were obtained by using a biprism which was inserted in front of the streak camera (4.18). By rotating the biprism with respect to the slit of the streak camera either the same or different regions of the plasma were imaged. If the same region was viewed, similar spectra were obtained. However, the spectra did not have the same temporal behaviour if two regions of the plasma separated by only 10 μm were imaged.

The backscattered spectrum and the burn through spatial image from a thin foil target were also recorded simultaneously as shown in Figure 4.06 where the average irradiance on target was $2 \times 10^{14} \text{ W cm}^{-2}$. In this set-up the backscattered signal was measured to be 45 ps earlier in arrival at the streak camera than the transmitted signal. Consequently the large reduction in the backscattered light (greater than the dynamic range of the streak camera) observed between 15 and 55 ps on the streak record corresponds to between 60 and 100 ps after the burn through first

appears. Using a self-similar planar expansion model to describe the motion of the foil we can calculate the sound speed from the foil thickness and the time delay from the beginning of the laser pulse until burn through. Assuming that the peak density in the foil at the time of burn through is the critical density n_c we can then calculate the peak density at the time when the backscatter turns off to be about $0.75 n_c$. We conclude that most of the backscattered light originates from densities between $0.75 n_c$ and n_c .

The multiplicity of small spots observed in the burn through images suggest that burn through occurs first in the intense regions of self-focused filaments. These filaments have been predicted theoretically and have recently been observed in laser produced plasmas. The new and important result of these observations is however that these filaments appear to be unstable in the critical density region with a lifetime of about 10 ps. Various models predict that the filaments may be unstable. A possible 'kink' mode similar to the MHD mode of pinched plasmas was suggested by Valeo (4.14). Parametric processes occurring within the filaments can produce strong ponderomotive pressures due to decay waves and thus destroy the filament. The locally increased light intensity can produce a higher ablation rate and destroy the filament by injecting more ablated material. It is also probable that filaments are surrounded by self-induced thermoelectric magnetic fields. The neighbouring filaments will attract each other like parallel currents. As an example a filament of 1 μm radius and surrounded by a magnetic field of 3 MGauss carries a current of 1.5 kAmp. Two such filaments 50 μm apart in a plasma of $n_e = 10^{21} \text{ cm}^{-3}$ will initially accelerate toward each other at a rate of $4 \times 10^{17} \text{ cm s}^{-2}$ and will coalesce in rather less than 50 ps. Clearly there are many mechanisms that could give rise to unstable filamentation but from the present data it is not possible to distinguish between them.

The unstable filamentation could also be a cause of the fluctuations in backscatter but we believe this to be unlikely since the fluctuations are seen from the beginning of the backscatter signal. The growth rates of filamentation and Brillouin scatter differ by a factor (v/ω_{1a}) , the

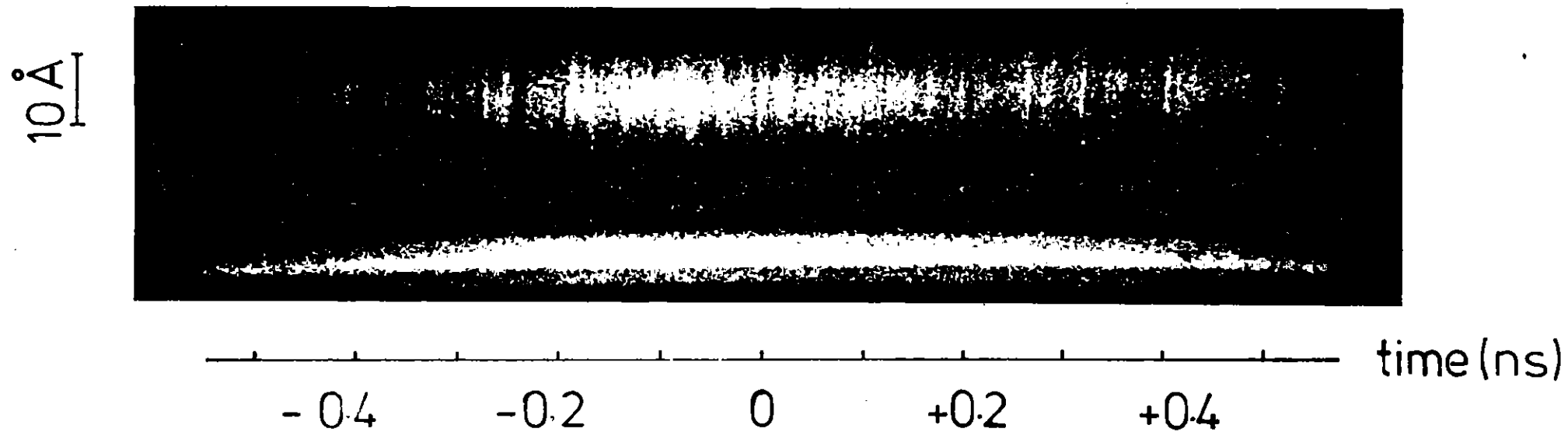


Fig. 4.05 Upper trace-spectrally resolved backscattered ω_0 signal.
Lower trace-temporal profile of the incident laser pulse.

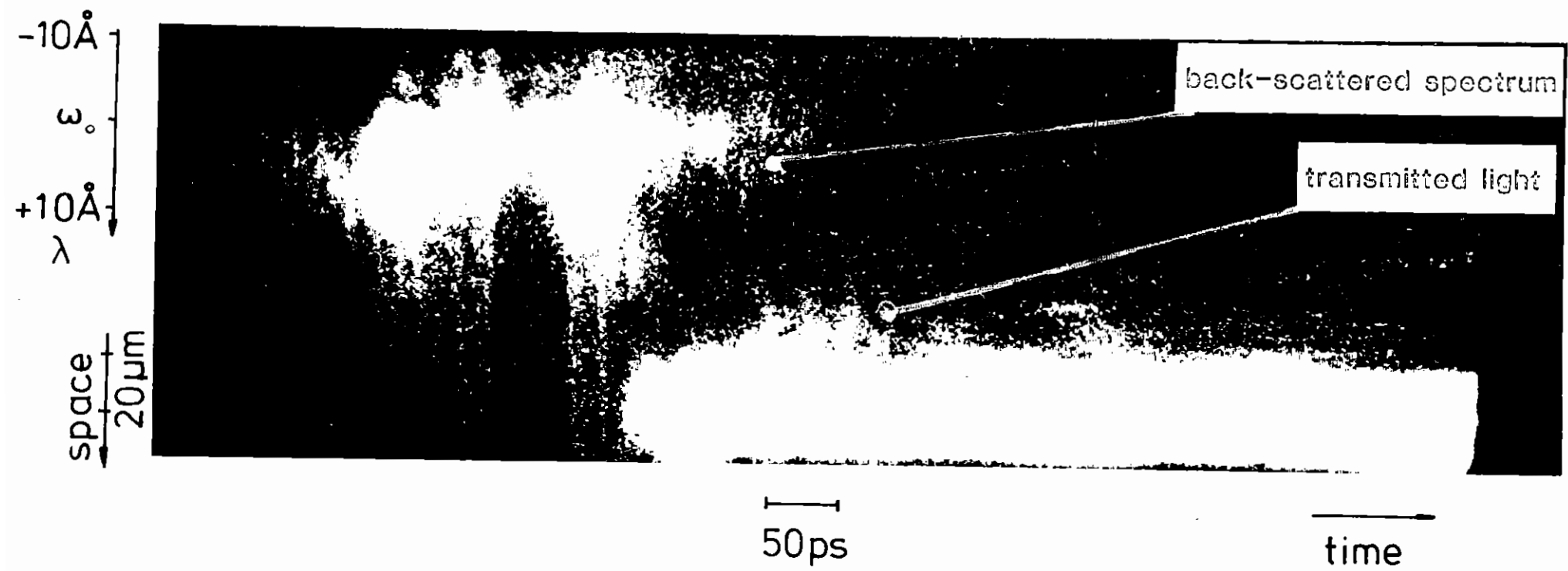


Fig. 4.06 Simultaneous backscatter and burnthrough signals on 0.5 μm foil target.

ratio of the ion wave damping rate to its frequency. For a hot plasma ($T_e \approx 300$ eV) and $ZT_e/T_i \gg 1$ the collisional and ion Landau damping are both small and the threshold for SBS will be exceeded long before there is sufficient under dense plasma to cause filamentation. However there is a straightforward explanation of the backscatter fluctuations purely in terms of SBS in an inhomogeneous plasma.

The ion wave driven by Brillouin backscatter always has a wave number k of about $2k_0$ where k_0 is the local value of the pump wave number.

$$k_0 = \frac{\omega_0}{c} \sqrt{1 - n_e/n_c}$$
Thus the driven ion wave frequency changes with density which is an expression of the limited region of the plasma over which any particular ion wave is amplified. It follows that the scattered electromagnetic waves also have slightly different frequencies and will interfere with each other with a characteristic time scale of $\tau = 1/\Delta\omega$. Since the behaviour of the pump wave near n_c will be described by an Airy function we will not actually have the ion acoustic waves extending down to zero frequency (ie $k = 0$) and so this time scale τ will be a few ion wave periods depending on the details of the plasma density profile. Using the value of the sound speed deduced earlier ($c_s \approx 10^7$ cm⁻¹) gives $\tau \approx 10^{-11}$ s, in very good agreement with the observations. The fluctuations are not truly periodic since the density profile evolves during the laser pulse.

In conclusion we have observed rapid changes in the spatial pattern of thin foil burn through which is most easily explained by unstable filamentation (or self focusing) of the laser light in the plasma. Rapid modulation of the backscatter light is probably caused by interference of SBS from different densities but may also be influenced by the unstable filaments.

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4.4 Filamentation Instability in Laser Produced Plasmas

Laser beam filamentation (where the laser beam breaks up into small

intense filaments) may play a vital role in laser plasma interaction and ablative implosion experiments. Since the filamentation instability is strongly dependent on the scalelength of the underdense plasma ie it is more effective in large, underdense plasmas, it gains in importance in ablative implosion experiments where long laser pulses (several ns) and large focal spots (up to several mm) are used. Initial small perturbations on the laser beams or on the target surface may grow due to self-focusing effects in the plasma. Various mechanisms, for example ponderomotive force, thermal expansion and relativistic effects can give rise to density depletions in the plasma leading to a focusing of either the whole or part of the laser beam due to the changes in the refractive index. The ponderomotive force processes are more likely to occur in collisionless plasmas whereas thermal effects may initiate self-focusing in collisional plasmas. Relativistic effects where the refractive index changes due to the mass change should be negligible in the experimental observations described below since they were carried out at irradiances of $I(\lambda/1 \mu\text{m})^2 \approx 10^{12} - 10^{15}$ W cm⁻².

The filamentation instability may be detrimental in various ways for laser driven implosions and may be vital for any target design. Since the irradiance is increased locally the threshold for resonance absorption and parametric instabilities may be exceeded locally resulting in the production of fast electrons and hard X-rays. In particular electrons may be trapped by the electron plasma waves which may originate from the two plasmon decay and stimulated Raman scattering instabilities and may be accelerated up to energies of several hundred keVs or even MeVs causing the fuel to be preheated. At shorter laser wavelengths the distance between absorption and ablation regions is small and there is little thermal smoothing of the non uniform illumination resulting from filamentation. The non uniform ablation pressure is then imprinted on the shell and an asymmetric implosion results.

The subcritical density region is investigated using various optical probing techniques such as interferometry, Faraday rotation, shadowgraphy and Schlieren imaging. A short pulse (70-100 ps) generated by a mode-locked oscillator which is synchronised to a long (0.5 to several ns)

pulse provides a short optical probing pulse. In order to study the time evolution of various physical processes such as laser beam filamentation, plasma jetting etc a double pulse optical probing system has been developed. It allows two images separated in time to be taken on each shot. The short gaussian pulse is amplified in a high repetition rate amplifier and is then passed through a multipass delay line in order to achieve synchronization with the main beam. The $1.052\ \mu\text{m}$ light is frequency doubled in a CDA crystal cut for noncritical phase matching and Raman shifted to the red ($\lambda = 622\ \text{nm}$) in a cell of ethanol for the purpose of frequency discrimination against the harmonics of the heating laser beams (4.11).

The $\sim 20\ \text{ps}$ red pulse is then split into orthogonally polarized pulses using a Wollaston polarizer. These are time delayed relative to each other by adjustable path lengths and after being recombined are directed past the target. The target is imaged onto film by an $f/3.3$ microscope objective with an overall magnification of 32. The two orthogonally polarized pulses are separated again after the target with a polarizing prism giving two images at different times. The polarizing prism is protected from the laser radiation scattered off the target by a red filter when the target is heated with green light or by a HA3 filter when the target is irradiated with infrared laser light. Interference filters ($100\ \text{\AA}$ in bandwidth) centred at $620\ \text{nm}$ are mounted on the camera to block radiation emitted from the plasma. The spatial resolution of the imaging system is 2 to $3\ \mu\text{m}$ in the object plane.

X-ray images are taken using a pinhole camera equipped with two pinholes $5\ \mu\text{m}$ and $12.5\ \mu\text{m}$ in diameter in a $10\ \mu\text{m}$ thick platinum substrate which are filtered with $15\ \mu\text{m}$ and $40\ \mu\text{m}$ of Beryllium respectively giving an e^{-1} transmission at $h\nu = 1.2\ \text{keV}$ and $1.6\ \text{keV}$ respectively. The pinholes are placed $2.5\ \text{cm}$ from the plasma and with an object to film distance of $10\ \text{cm}$ a $4\times$ magnification on the film is obtained. A spatial resolution of $6.5\ \mu\text{m}$ is achieved with the $5\ \mu\text{m}$ pinhole. The images are recorded either on Kodak No-Screen film or Kodak Industrex film.

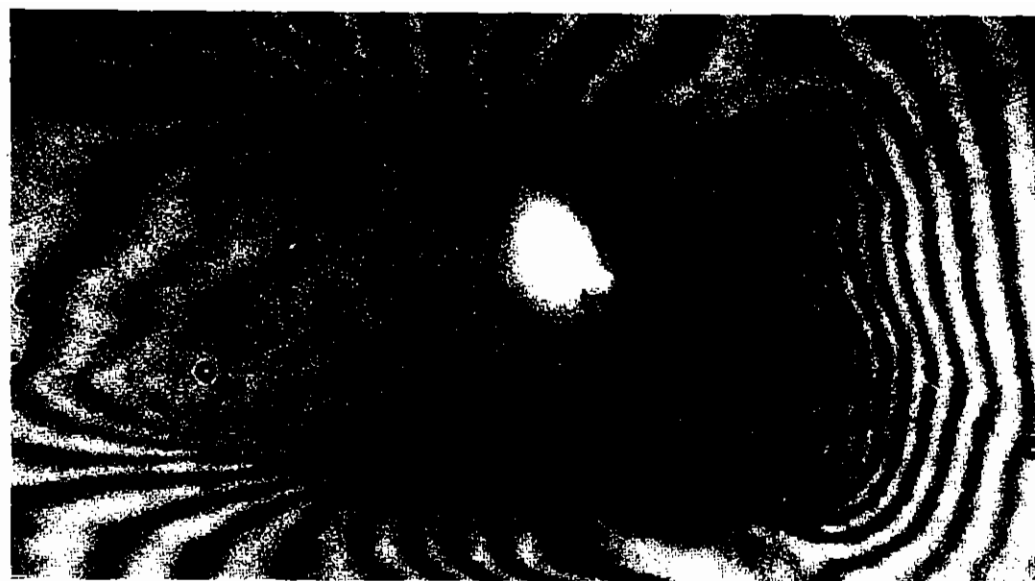
Laser fusion targets are much larger (several mm in diameter) than the

targets (a few $100\ \mu\text{m}$ in diameter) used for previous experiments and consequently are surrounded by much larger scalelength plasmas. In order to simulate this situation longer scalelength plasmas were produced by irradiating thin foil targets with a large prepulse.

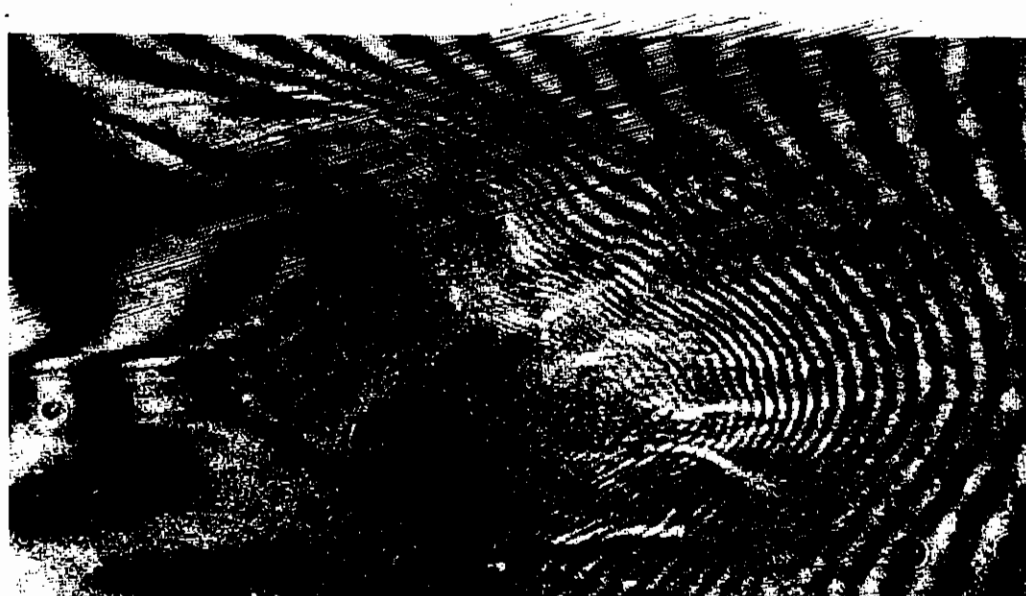
Figure 4.07(a) shows an interferogram recorded on a plastic disc $270\ \mu\text{m}$ in diameter and $2\ \mu\text{m}$ in thickness. The target was irradiated with a single $100\ \text{ps}$ (FWHM) infrared pulse at an intensity of $1 \times 10^{15}\ \text{W cm}^{-2}$. The interferogram was taken $4\ \text{ns}$ after the peak of the laser pulse using a $350\ \text{nm}$ probing beam. Notice the planar expansion of the plasma. Figure 4.07(b) shows an interferogram recorded on a similar target which was however irradiated with two pulses separated by $4\ \text{ns}$. The interferogram was recorded shortly after the peak of the second pulse. 60% of the incident laser energy was concentrated in the first pulse and with a focal spot of $100\ \mu\text{m}$ in diameter the target was irradiated at an intensity of $6.7 \times 10^{14}\ \text{W cm}^{-2}$. Since the preformed plasma expanded up to $\sim 200\ \mu\text{m}$ towards the laser beam (which can be seen in Figure 4.07(a)) and an $f/1$ lens was used to focus the laser beam, the second pulse was focused to a focal spot of $\sim 300\ \mu\text{m}$ in diameter giving an estimated irradiance of $\sim 5 \times 10^{13}\ \text{W cm}^{-2}$. Filamentary structures can clearly be seen in Figure 4.07(b) where the laser beam is incident. In addition however a thin spike about $28\ \mu\text{m}$ in diameter is observed on the rear side of the target suggesting that the whole beam may have self-focused.

Filamentation may cause the laser energy to be deposited locally and consequently in order to obtain a uniform ablation pressure to drive an implosion any non-uniformities have to be washed out between the absorption and ablation region. However, at shorter laser wavelengths thermal smoothing effects become inefficient arising from the short distance between the absorption and ablation region.

Experimental observations were carried out to investigate the effects of filamentation on the ablation pressure on planar targets irradiated with a UV laser beam where tiny spots were clearly observed in X-ray pinhole camera images. The laser beam was focused onto the target using an $f/2.5$



← LASER



50 μm

Fig. 4.07 a) Interferogram taken a 2 μm thickness plastic disk, irradiated with a single laser pulse.
 b) As a) but irradiated by two laser pulses 4ns apart.

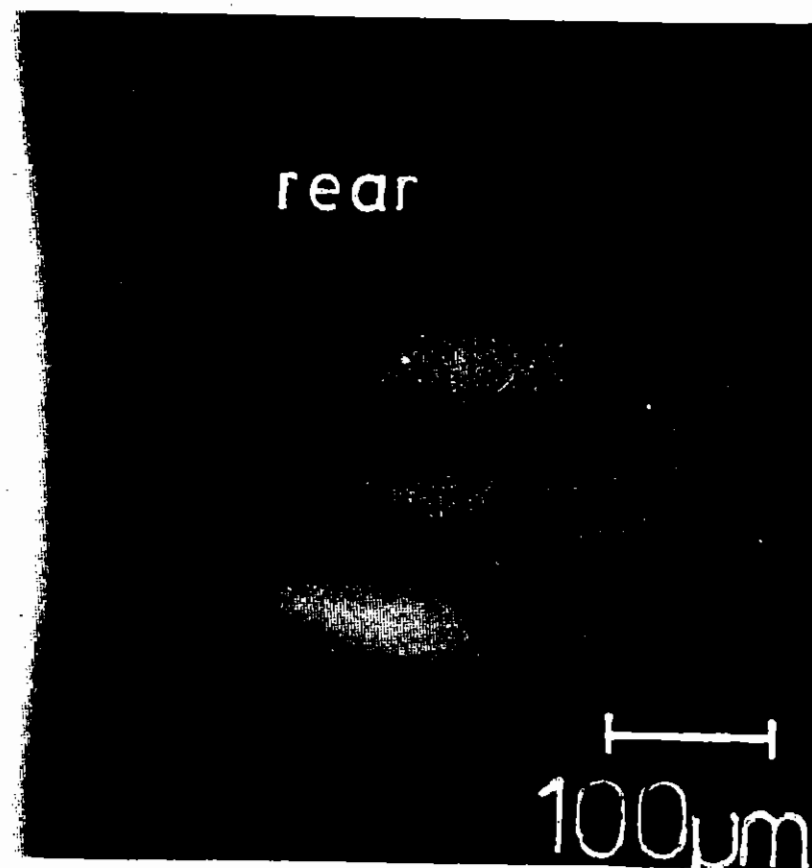
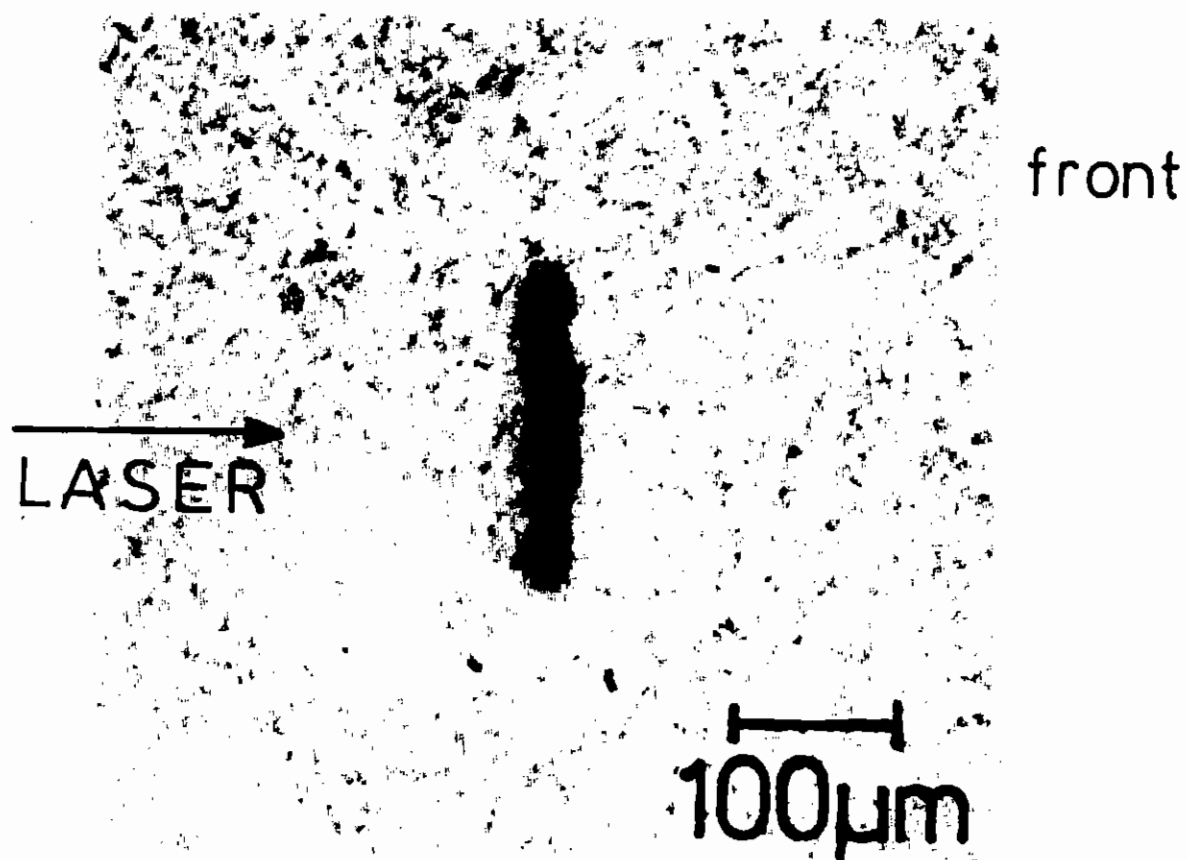


Fig. 4.08 X-ray pinhole camera ($\sim 1\text{keV}$ energy) and optical images showing non uniform shock breakthrough caused by laser beam filamentation.

fused silica spherical doublet lens and when the target was displaced from the best focus of the lens a ring structure appeared on the target surface arising from the spherical aberration in the lens. On medium and high Z targets the annular ring broke up periodically into tiny spots. Observations were carried out at 90° to the direction of the incident beam using an X-ray pinhole camera and an f/3.3 microscope objective which imaged the plasma onto a camera mounted with an interference filter centred at 456 nm (5 nm bandwidth) with a 30 X magnification. Figure 4.08 shows an X-ray pinhole camera image and a time integrated optical image taken on a gold foil 4 µm in thickness which was irradiated with a 0.35 µm laser beam. The irradiance on target was $3.2 \times 10^{13} \text{ W cm}^{-2}$. Tiny spots can clearly be seen in the X-ray image. On the rear side of the target foil structures with the same spatial scalelength as the X-ray spots can be seen in the optical image. This indicates that non uniform shock break-through occurs if spots in the X-ray pinhole camera image are observed. Similar observations were also made on Al foils 25 µm in thickness. The optical self-emission was however more uniform on plastic foils which is consistent with X-ray images taken on plastic targets where more uniform X-ray images were obtained.

The regime in which no non-uniformities in the ablation pressure may be expected is shown in Figure 4.09 where the dependence of the thermal smoothing and the self-focusing effects on the irradiance and laser wavelength are plotted. The scaling law for the ponderomotive force self-focusing was derived by Max and Estabrook (4.19) and is given by

$$I \geq 2 \times 10^{15} \frac{T_e}{L \lambda_L (n_e/n_c)}$$

where I is the incident irradiance in W cm^{-2} , T_e is the electron temperature in keV, L is the density scalelength in microns, λ_L is the laser wavelength in microns, and n_e and n_c are the electron density and critical density respectively. On the other hand a wavelength scaling for thermal diffusion effects depending on the separation between the absorption region and ablation front was obtained by numerical calculations (4.20). Further analysis gives (4.21)

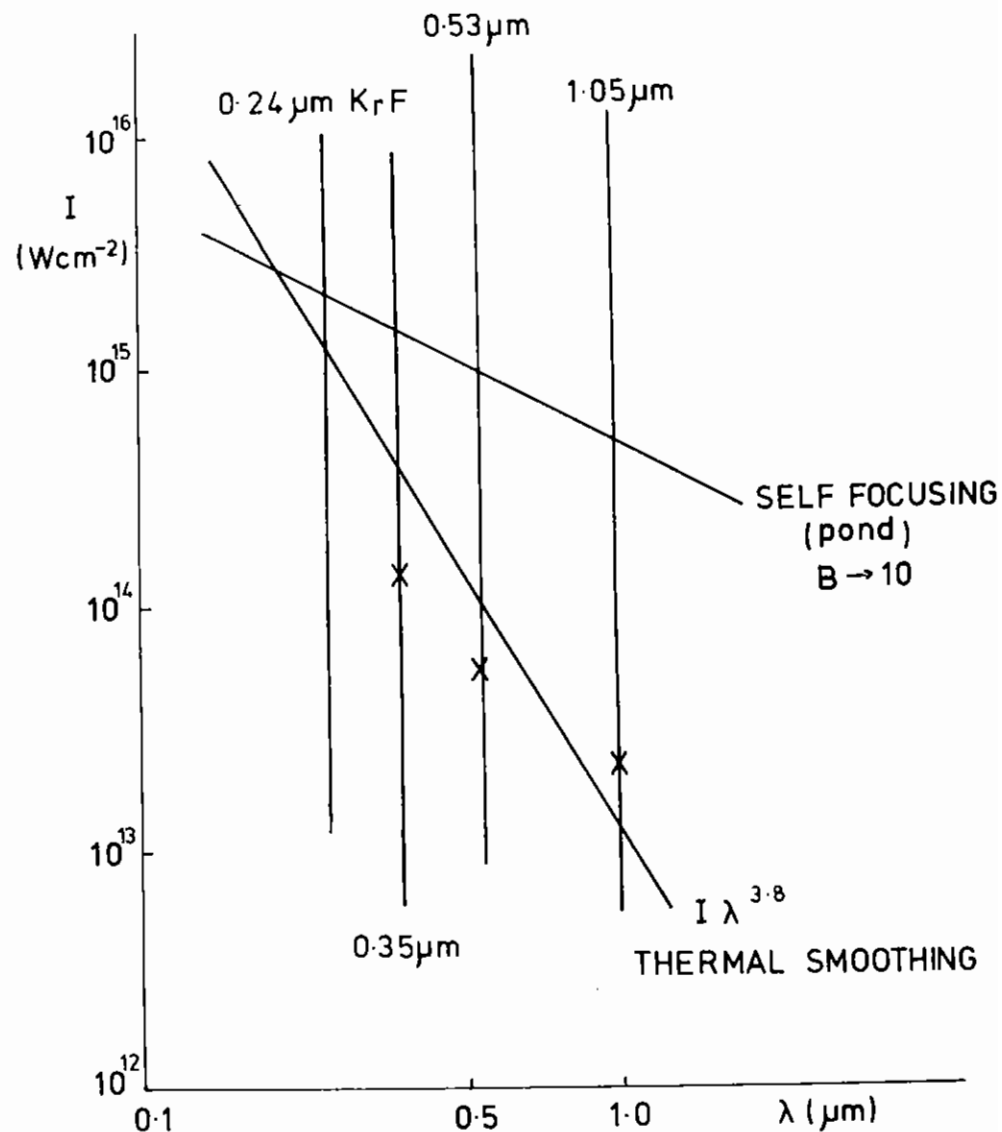


Fig. 4.09 Irradiance and wavelength scaling for thermal smoothing and self-focusing effects. The crosses represent experimental observations which clearly showed filaments.

$$I \geq 3 \times 10^{13} \lambda_L^{-3.8}$$

where I is again the irradiance in W cm^{-2} and λ_L is the laser wavelength in microns. To obtain a uniform ablation pressure the experimental parameters have to lie between the two curves. The crosses in Figure 4.09 represent measurements made on plastic targets at three different laser wavelengths using optical probing techniques. At the indicated irradiances filamentary structures aligned along the laser beams were clearly observed. These values of irradiance are well below the theoretically predicted curve which is probably caused by initial perturbations on the laser beams. The crosses at shorter laser wavelength are below the thermal smoothing curve and consequently non-uniformities in the ablation pressure may exist at these irradiances which is consistent with X-ray pinhole images where some spottiness was observed. Notice that non-uniformities in X-ray pinhole camera images were observed even at lower irradiances than those plotted in Figure 4.09 when high Z-targets were used. Thermal self-focusing effects which are not shown in Figure 4.08 may however have been responsible for the filamentary structures in those observations.

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CHAPTER 5 TRANSPORT AND PARTICLE EMISSION STUDIES

INDEX

- 5.1 INTRODUCTION page 5.1
- 5.2 FAST ELECTRONS IN SPHERICAL GEOMETRY page 5.3
- 5.3 MEASUREMENT OF K-EDGE BROADENING CAUSED BY PREHEAT FROM LASER PLASMAS page 5.4
- 5.4 ABSORPTION SPECTRA FROM 1-2 TRANSITIONS IN HEATED C₂ AND Si AT SOLID DENSITIES page 5.10
- 5.5 X-RAY SPECTROSCOPY OF LOW OPACITY POLYMER TARGETS page 5.11
- 5.6 THERMAL SMOOTHING EFFECTS UNDER STEADY-STATE CONDITIONS page 5.15
 - 5.6.1 Experimental Investigation of Long Pulse Smoothing
 - 5.6.2 Results
 - 5.6.3 Simulations of Long Pulse Smoothing
- 5.7 THERMAL SMOOTHING MEASUREMENTS BY SHOCK TRANSIT TIME page 5.20
- 5.8 A MEASUREMENT OF THE STOPPING POWER OF HOT DENSE MATTER page 5.24
- 5.9 LASER SCATTERING MEASUREMENTS OF HEAT FLOW AT HIGH TEMPERATURE GRADIENTS page 5.31

REFERENCES page 5.36

CHAPTER EDITOR: W T Toner

CHAPTER 5 TRANSPORT AND PARTICLE EMISSION STUDIES

5.1 Introduction

The activities of the Transport and particle Emission Group in 1982 covered a wider field than usual. Familiar techniques were applied in new areas, in particular in dense matter physics, and additional techniques were applied to old problems such as thermal smoothing and transport inhibition.

It has been suggested in earlier reports that fast electron preheat in long pulse irradiation might be very much less than observed with short pulses and it has been shown that the bulk of long pulse preheat in planar geometry occurs at large distances from the focal spot. In section 5.2 data are presented which show very low levels of long pulse preheat in spherical geometry indicating that very little of the fast electron energy couples to the target. What little preheat was seen could be due to radiation.

Two very different aspects of the physics of hot dense matter were the subject of experiments: broadening of features in X-ray absorption spectra is described in 5.3 and 5.4 and changes in alpha particle stopping power in 5.8. The purpose of the time-resolved X-ray absorption experiment described in 5.3 - the first such measurement - was to observe an anticipated shift in the K edge of an ionised medium due to ionisation depression. No shift was observed but the edge was substantially broadened. Analysis shows that the ionisation depression is almost exactly cancelled by an increase in inner shell ionisation energy due to outer shell ionisation and that the edge broadening can be accounted for by statistical variations in the continuum depression. The use of edge broadening as a sensitive preheat diagnostic is discussed. In section 5.4 a preliminary analysis of shifted and broadened absorption line spectra from shock compressed and heated solid absorbers is presented.

The possibility of ion beam fusion has stimulated interest in the ion stopping power of hot dense plasmas. A novel experiment which might be

said to be the first alpha particle backlighting experiment is described in section 5.8. Alpha particles from an exploding pusher implosion were used as test particles whose energy loss in a separate, laser heated, mylar foil was measured. In a first trial of the method the stopping power of the hot foil was found to be a factor of 1.33 times greater than that of the same foil when cold, in agreement with theoretical estimates.

A new type of gas-filled target with a polymer shell is described in 5.5. The advantages resulting from the low X-ray opacity and emissivity of these shells are demonstrated by a study of chlorine emission spectra from the cores of exploding pusher implosions. Such targets are well suited to emission spectroscopic diagnosis of ablative compression implosions.

Further work has been done on thermal smoothing of lateral variations of ablation pressure and its wavelength and irradiance dependence. The X-ray backlighting measurements reported last year were improved on by the use of a longer drive pulse to reduce the effect of foil decompression. The wavelength scaling between 1.05 μm and 0.53 μm agrees with simulation results but the irradiance dependence is stronger than predicted. Some decompression effects were still evident. Section 5.7 contains preliminary results of an experiment using a shock transit method to measure lateral variations in ablation pressure which is capable of being used at early times when decompression effects should be less marked. This is the first experiment to observe smoothing at 0.351 nm.

The question of transport inhibition in large temperature gradients is still open. A trial was made of a laser scattering method developed elsewhere by the Imperial College group to measure thermal conductivity in a laser heated, low density, Z-pinch plasma. An optical streak camera was used to time resolve the spectrum of the scattered light with high resolution, for the first time. It was demonstrated that the large values of ∇T needed to attack the transport inhibition problem could be generated. Strong enhancement of the ion feature was observed, indicating the presence of ion turbulence and an interesting asymmetry in the

spectrum of the electron feature was seen which may indicate substantial electron flow.

W T Toner

5.2 Fast Electrons in Spherical Geometry

Planar geometry experiments (5.01)(5.02) with long (> 1 ns), $1.05 \mu\text{m}$ laser pulses have shown that although the fast electron energy deposition directly beneath the focus is small (typically less than $\frac{1}{2}\%$ of incident energy for 1.3 ns pulses) there is significant ($> 5\%$ of incident) deposition of energy by fast electrons at distances of up to 5 mm from the focal spot in the plane of the target. In spherical geometry there is no possibility of radial transport, and this "remote energy deposition" could significantly increase the bulk preheat if it couples efficiently to the target.

The aim of this experiment was to determine the fast electron energy deposition for spherical irradiation at $1.05 \mu\text{m}$ with 1.5 ns pulses. This was to be achieved by measuring the K_{α} emission from solid vanadium spherical targets with various thicknesses of coatings, using both time integrating and streaked X-ray spectrometers. Unfortunately, only a short run was possible. However, the four useful data points we obtained showed that the planar geometry remote deposition does not couple to the target in spherical geometry and the preheating is small.

The experimental set up used 6-beam irradiation of solid V spheres, with 150 - 300 J of $1.05 \mu\text{m}$ energy incident in 1.5 ns. The targets used were 60 - 220 μm in diameter, either bare or with a 2 - 6 μm parylene coating. Incident intensities were in the range $(1.5 - 15) \times 10^{14} \text{ W cm}^{-2}$.

The spectra were recorded on a miniature time integrating spectrograph using either a Ge(2d = $4\text{\AA}$) or LiF (2d = $4\text{\AA}$) Bragg reflecting crystal, and on the streaked X-ray spectrometer using the Ge crystal. The time integrated data is shown in Table 5.01.

Table 5.01 Time integrated data for fast electron energy deposition in spherical targets

Shot	Target diameter in μm	Intensity (W/cm^2)	K_{α} (J/Sr)
14210482	66	14.5×10^{14}	8.6×10^{-6}
08160482	89	6.1×10^{14}	7.0×10^{-6}
09070482	71	5.9×10^{14}	4.8×10^{-6}
11070482	80	5.4×10^{14}	3.9×10^{-6}
02070482	89	5.0×10^{14}	4.2×10^{-6}
15070482	97	3.6×10^{14}	4.9×10^{-6}
11060482	107	2.8×10^{14}	1.4×10^{-6}
03130482	71 + 2 μm CH	6.6×10^{14}	> 0

In all shots the $V K_{\alpha}$ yield was small, the largest measured yield being $8.6 \times 10^{-6} \text{ J Sr}^{-1}$. It is thought that this yield may be dominantly radiation pumped (the $V XXII 1s3p - 1s^2$ can efficiently excite $V K_{\alpha}$ emission). However, taking the measured K_{α} yield as an upper limit to the fast electron induced yields gives an energy deposition of 0.1 J out of 140 J incident (5.03). Problems with high laser prepulse levels made the data from targets with plastic overlays difficult to interpret, though on the one shot where the thin plastic layer suppressed the V plasma emission the $V K_{\alpha}$ was not seen. This also tends to suggest that radiation is the dominant pump for K_{α} emission, implying that the fast electron position may be significantly below the upper limit of 0.07 % estimated above.

The single streak record with K_{α} emission shows that the emission follows the level of the plasma X-ray continuum as would be expected with photopumping.

The experiment is to be repeated later this year in collaboration with colleagues from LASL. It is important to measure the changing behaviour of fast ion emission and preheat with target size, fast ions being the likely sink for the radially spread fast electron energy in planar experiments. The data already taken does show that the deposition is small - it is possible that for any optimised ablative implosion, the preheating will be dominantly due to X-rays.

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5.3 Measurement of K-Edge Broadening caused by Preheat from Laser Plasmas

The width of the K shell photoabsorption edge in chlorine heated by laser plasmas is seen to vary during the laser pulse. A broadening but negligible shift of the edge, which correlates with the energy deposition is observed. It is argued that this is due to microfield variations in the continuum depression.

It is well known that when high intensity lasers interact with solids a

high temperature plasma at about the critical density for the wavelength of the incident laser is formed. However at a depth of a few μm into the solid a much cooler plasma, albeit at approximately solid density is formed, by the shock wave propagating through the solid (5.04), or by radiative (5.05) or fast electron (5.06) preheat. This plasma is strongly coupled, and can be degenerate. However its high density and low temperature (10 eV - 50 eV) necessitates the use of absorption rather than emission spectroscopy for diagnosis.

X-ray absorption spectra have been taken before (5.07 to 5.10) from laser produced plasmas using the continuum from the highly emitting core of compression targets to backlight the surrounding glass which has been heated by fast electrons, resulting in absorption lines from helium and lithium like states of ionisation.

Here we report the first time-resolved measurements of X-ray absorption spectra of solids heated to temperatures of $\sim 10 \text{ eV}$ by X-rays from laser produced plasmas.

The experiments were motivated by the expectation of a spectral shift of the K shell photoabsorption edge (of chlorine): no shift was observed but the edge was seen to become broader during the laser pulse. A correlation was observed between the time integrated width of the edge and the independently measured energy deposition by X-ray preheat in an absorber layer. The absence of a shift is explained as a result of an increase in the photoionisation energy due to ionisation cancelling the decrease due to ionisation depression. The width of the edge is due to a statistical variation of the continuum depression.

For targets with the absorber layer close to the X-ray emitter the width of the edge was used to infer the degree of preheat.

The spectral position of the photoionisation edge of an isolated atom shifts to higher energy as the atom becomes ionised, because of reduced screening. Atomic calculations (5.10) show that the increase in the inner shell photoionisation energy is essentially equal to increase of the outer

shell ionisation potential. If the ion is immersed in a plasma there is a depression of the continuum (5.12) which, for the case of the ion sphere radius $r_s = (3/4\pi Na)^{1/3}$ being much less than the Debye length, is $\Delta E_K = 3Ze^2/2r_s$. Estimates of these two effects for chlorine in both solid density KCl and shock compressed KCl (at 10 Mbar) are shown in Figure 5.01 as a function of the state of ionisation of the solid. As can be seen the two effects in Figure 5.01 nearly cancel. Furthermore there is some uncertainty in the value of the continuum depression for strongly coupled plasma (5.13) hence it is unclear whether the K edge would shift to higher or lower energy. Additionally, as discussed below, a broadening of the edge might be expected.

Experimental measurements of the position of the K edge of Cl in KCl were made using a 0.20 J, 0.53 μm wavelength laser beam with pulse lengths of 0.4 ns and 0.8 ns. The beam was focused by a f/1 lens. The irradiance was varied by changing the focal position of the target, giving a minimum spot size of 35 μm and a corresponding maximum irradiance of $2 \times 10^{15} \text{ W cm}^{-2}$.

The targets had a layer of KCl (Cl absorber) buried at various depths (0 μm , 5 μm and 12 μm) beneath mylar and a surface layer of Bi. Bismuth was used as it has a bright X-ray emitting continuum near the Cl K edge when irradiated by the laser beam. This Bi laser produced plasma was viewed through the buried KCl absorber by a time resolving and a time integrating spectrometer at the back of the target as shown in Figure 5.02. Time resolved absorption spectra were recorded on a streaked X-ray crystal spectrograph with temporal and spectral resolution of 150 ns and 3 eV respectively. A time integrating spectrometer was also positioned at the front of the target to measure the spectral shape of Bi 'backlighting' emission in the region of Cl absorption.

Figure 5.03 shows a streak record of the Cl K edge. For this shot the laser irradiance was $1.4 \times 10^{15} \text{ W cm}^{-2}$ onto a target of 0.7 μm Bi, 5 μm KCl, 5 μm mylar. The microdensitometer tracings in Figure 5.03 show that initially the photoabsorption edge has a width of 10 eV which increases with time, reaching 40 eV by the end of the pulse. The mean position of

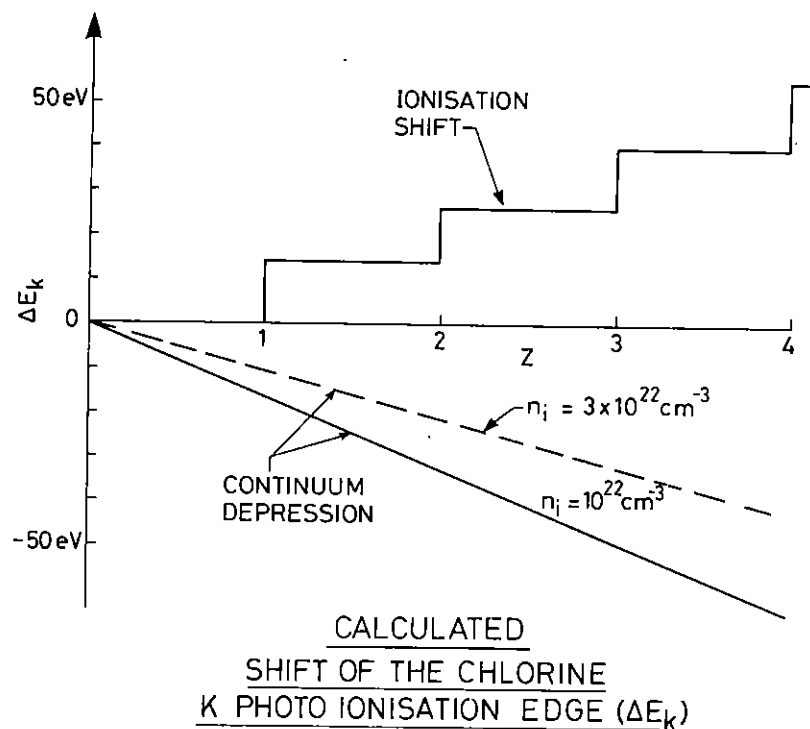


Fig. 5.01 Calculated changes in photoionisation energy of Cl in KCl vs state of ionisation for solid density KCl and shock compressed KCl.

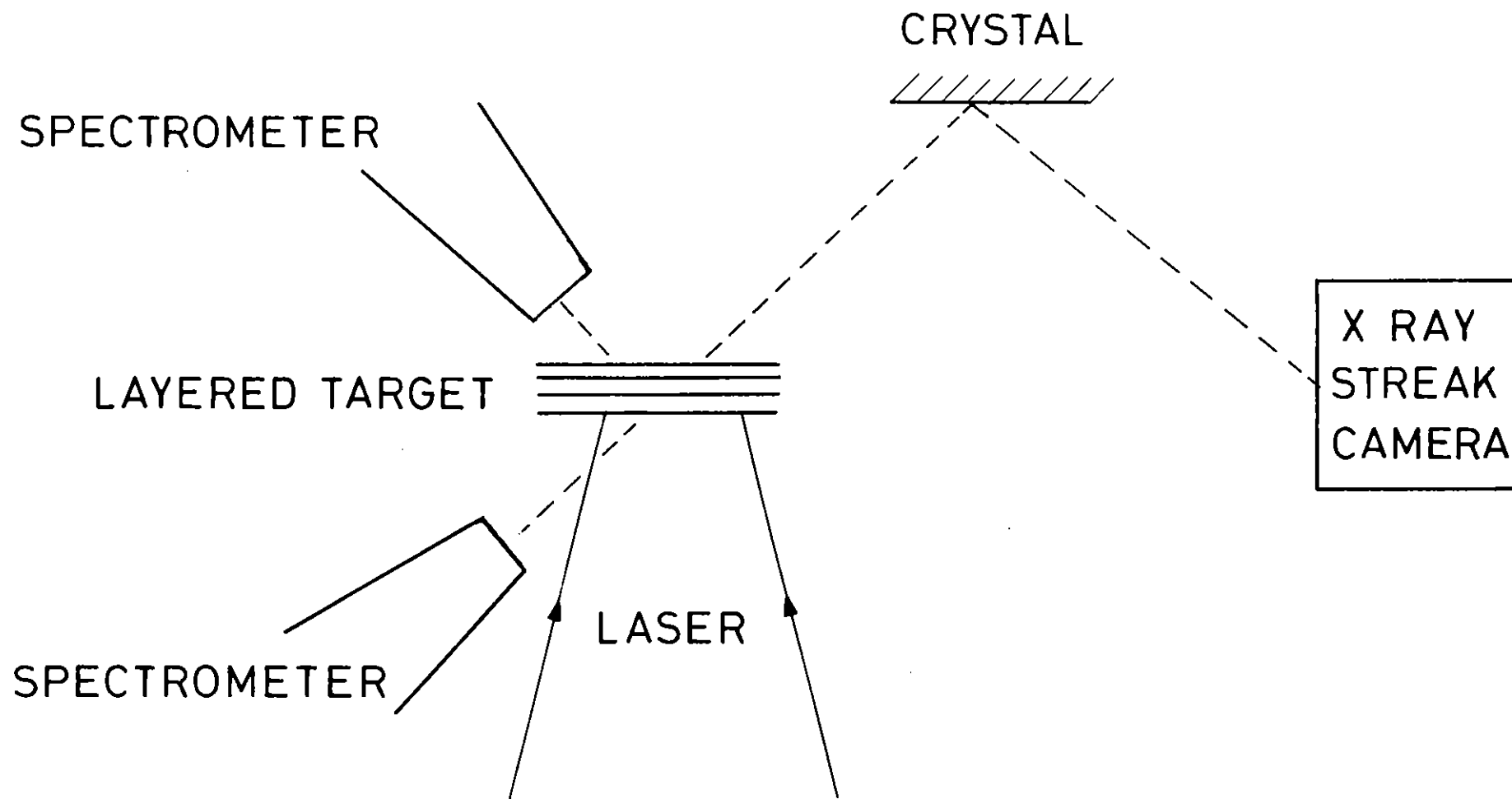


Fig. 5.02 Experimental arrangement for K edge experiment.

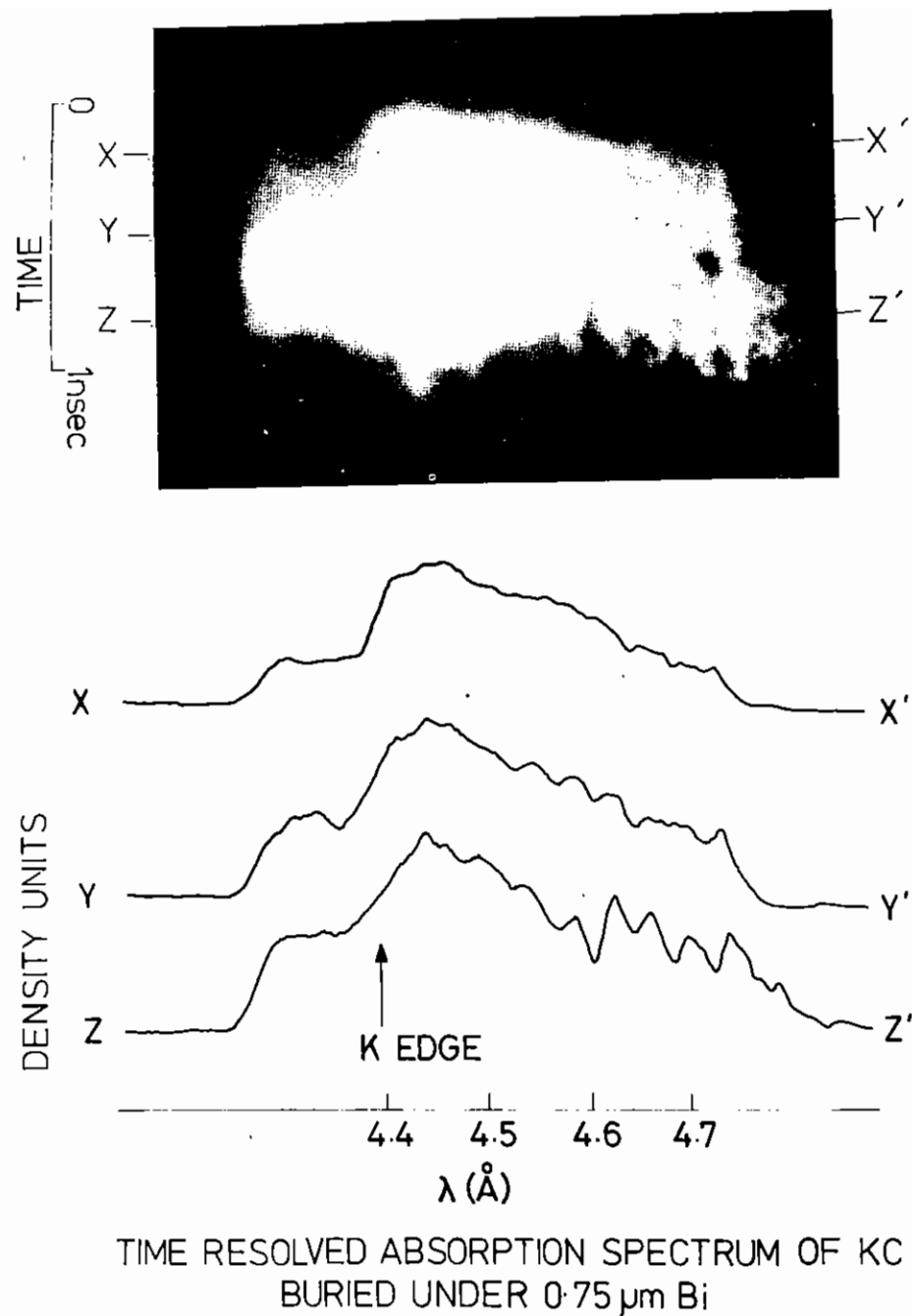


Fig. 5.03 Streaked Cl absorption spectrum at high intensity showing broadening of K edge during laser pulse.

the edge does not change within a measurement error of 5 eV, caused principally by the 'pin cushion' distortion of the streak-camera and intensifier.

Figure 5.03 also shows Cl XVI $1s^2 - 1s2p$ emission lines (at 4.45 Å) late in time when the Bi has ablated off. Absorption lines are also visible which have been identified as $1s-2p$ transitions in various ionisation stages of Cl (5.11). These absorption features were only observed when the KC1 layer was close to the front of the target ie just beneath Bi. Broadening of the K edge was however observed when the KC1 layer was buried under 12 μ m of mylar.

The degree of preheat of the absorber layer was varied by altering its depth in the substrate. During the laser pulse only a small amount of material was ablated from the surface (5.14). Hence the heating of the absorber layer was due to X-rays and shock waves from the Bi surface plasma, the fast electron reheating being negligible for the values of $I\lambda^2$ used here. The X-ray preheating was calculable and varied due to different absorption in the various thicknesses of mylar used. Figure 5.04 shows the Bi spectrum viewed from the front and two absorption spectra, one for a buried and therefore low temperature absorber, and one for a high temperature absorber corresponding to a layer of KC1 just below the Bi.

The low temperature absorber in Figure 5.04 shows the Cl K α emission at a high level, indicative of radiative and not fast electron K shell ionisation. The K α line acts as a fiducial and allows an accurate measurement of the position of the K edge to be made. There is no shift in the wavelength of the K edge but the edge width is 10 eV compared to a source width of 5 eV. The high temperature absorber in Figure 5.04 shows a much broader but unshifted edge.

Its absolute position was obtained by fitting a dispersion relation to the K α line and the Cl XIV emission line seen. This allowed the absorption features to be identified (5.15)(5.16) as $1s-2p$ transitions in ionisation stages of Cl from Cl IX to Cl XIV with Cl XVI appearing in emission.

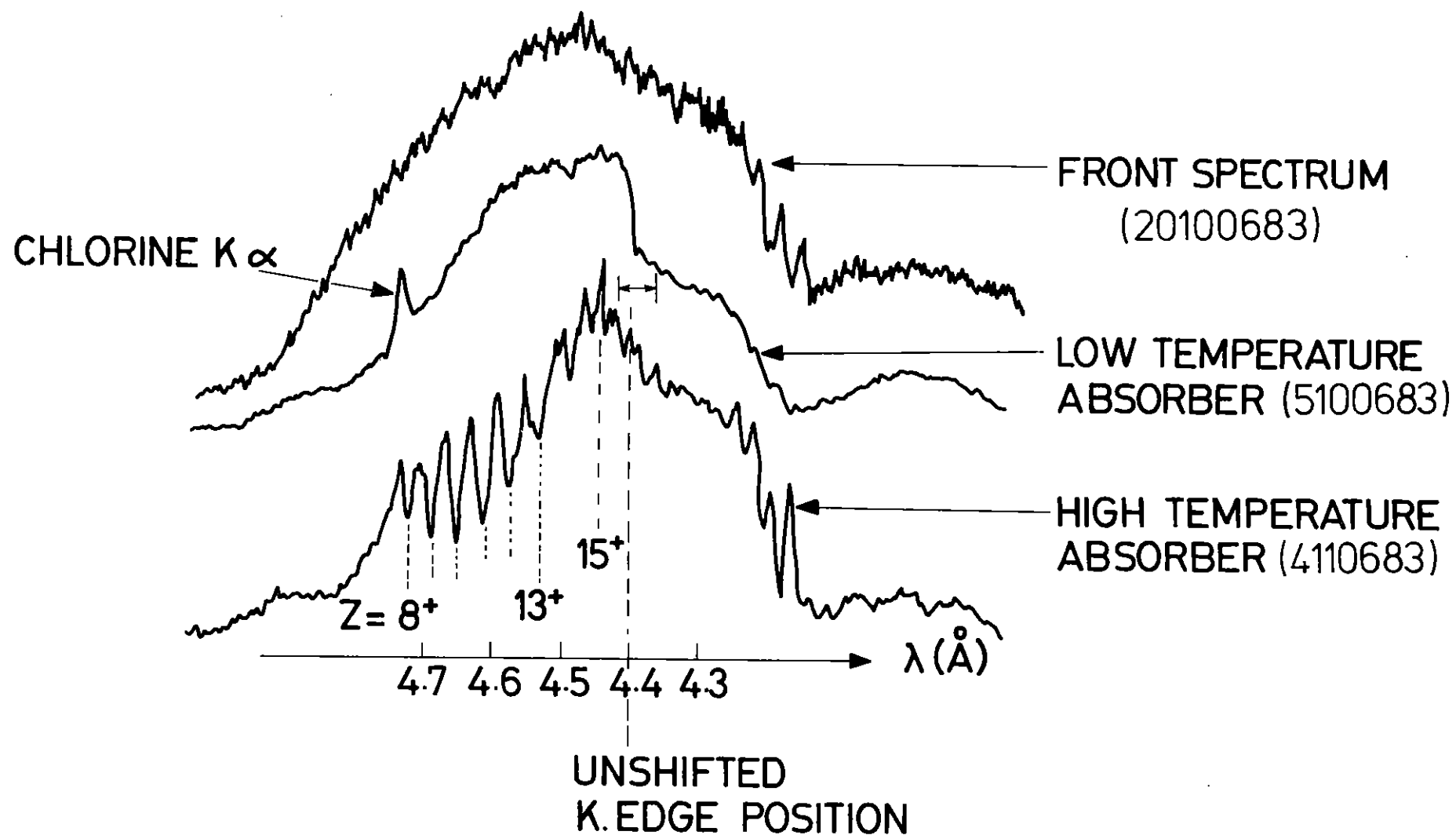


Fig. 5.04 Time integrated spectra of both low temperature and high temperature Cl absorbers, also showing shape of unfiltered backlighting spectrum.

Time resolved and time-integrated data were obtained for a wide range of targets and irradiation conditions. The intensity of K α emission was not a function of absorber depth. Fast electrons at these low values of $I\lambda^2$ have a relatively short range and therefore the K α emission is driven by photoionisation by photons of energy greater than the K shell photoionisation energy IP_K , which have a large range in Mylar.

Buried absorber layers are subject only to radiative preheat that is able to penetrate the Bi layer (until it is ablated) and the 5 μ m or 12 μ m Mylar. Radiation with $h\nu > IP_K$ will produce K shell ionisation and hence K α emission.

For a given absorbed spectrum $\frac{\partial E}{\partial h\nu}$ the K α emission is (5.10):

$$Y = h\nu K_\alpha \cdot \left(1 - \frac{1}{J_R}\right) \cdot \omega \int_{IP_K}^{\infty} \frac{\partial E}{\partial h\nu} \cdot \frac{1}{h\nu} \cdot dh\nu$$

and the absorbed energy is

$$E_{abs}(h\nu > IP_K) = \int_{IP_K}^{\infty} \frac{\partial E}{\partial h\nu} \cdot dh\nu$$

where J_R is the K absorption jump ratio and ω is the fluorescence yield.

The form of the absorbed spectrum in the region of the Cl K edge was obtained from the comparison of front and rear spectra for several shots and account was made of angular variations of absorption length to obtain $\partial E/\partial h\nu$. Thus Y was related to $E_{abs}(h\nu > IP_K)$, assuming the form of the absorbed spectrum was the same for all the targets. The size of the radiatively preheated volume was obtained from the X-ray pinhole image, the spatial extent of the K α emission region, the spatial extent of the continuum source and the thickness of the absorber. A preheat level in eV per atom by photons of $h\nu > IP_K$, $EP(IP_K)$, was thus obtained. Figure 5.05(a) shows a plot of average edge width from both static and streaked

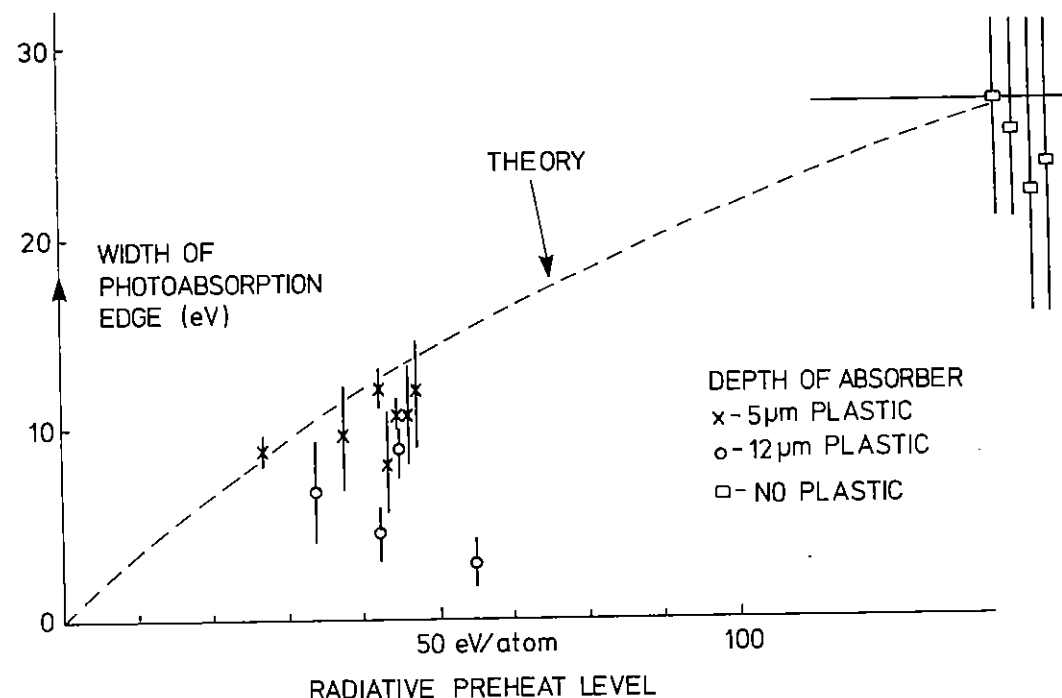
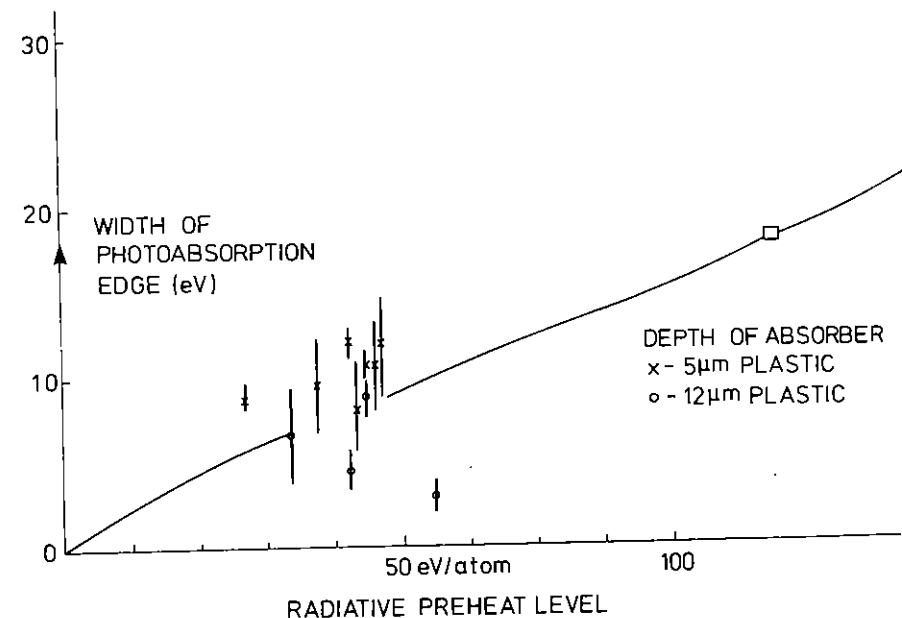


Fig. 5.05 a) Graph of measured edge width vs preheat level for KCl absorbers buried under 5 μ m and 10 μ m Mylar. Theoretical curve is also shown.
b) Edge width vs preheat level for Cl absorber directly underneath the Bi backlighting plasma.

spectra as a function of $EP(IP_K)$ for the three target types. Note that the broadening for the 5 μm and 12 μm Mylar targets is similar, suggesting that 7 μm Mylar does not significantly attenuate the preheating spectrum. We conclude that for these target types the radiative preheat is dominantly due to the spectrum in the 2 keV region which is recorded.

For targets with the absorber on the surface the broadening was much greater, as shown in Figure 5.05(b). For this configuration the X-ray preheating would not be due solely to X-rays in the region of the C & K shell. There would be a contribution from soft X-rays which would not be able to penetrate the mylar surface layer.

A simplistic model can be used to account for the broadening shown in Figure 5.05(a). The absorber layer is at a high density corresponding to a shock compressed solid. The ion density is about 10^{23} cm^{-3} and the temperature is of $\sim 20 \text{ eV}$, thus the plasma is strongly coupled. The continuum depression would therefore be given by the ion sphere model (5.13).

$$\Delta E_K = - \frac{3}{2} \frac{Ze^2}{r_s}$$

(where r_s is the radius of the ion sphere) as plotted in Figure 5.01. For a gas of uncorrelated particles the probability of the nearest neighbour being at a distance r is given by

$$W(r) = 3(r/r_s)^2 \exp(-(r/r_s)^2)$$

If the shift due to continuum depression is $\Delta E_K = - \frac{3}{2} \frac{Ze^2}{r}$ and is weighted by $W(r)$ then the width of the edge, $\delta(\Delta E_K)$, would be given by $\delta(\Delta E_K) = 0.76 \times \frac{3}{2} \frac{Ze^2}{r_s}$.

The dependence on specific energy arises from the Z dependence in this expression. Using an LTE approximation, the energy deposition of Figure 5.05(a) can be related to Z and hence an edge width corresponding to the

energy deposition. This is plotted in Figure 5.05(a) and is in broad agreement with the observed edge width for buried absorbers.

An extrapolation of this crude theory to the observed edge width for the laser of the surface absorber (of 20 eV) allows an estimate to be made of the energy deposition for this case of $\sim 150 \text{ eV/atom}$, as shown in Figure 5.05(b). Estimates of the X-ray conversion efficiency and hence soft X-ray preheating of a surface layer of the KC ℓ absorber lead to similar values of the preheating, although over the relatively thick absorber there would be a large variation in the level of heating.

In conclusion, we have made the first time resolved measurements of the C & K absorption edge in high density matter. The edge is observed to broaden in response to absorber heating in approximate agreement with a simplistic model based on statistical variations in the microfield. The degree of absorber radiative preheating is obtained by measuring the radiatively induced fluorescence in the absorber. The K edge width in material subject to unknown (probably soft X-ray) preheating is tentatively used as a preheat diagnostic.

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5.4 Absorption Spectra from 1-2 Transitions in Heated C ℓ and Si at Solid Densities

Targets were made unintentionally with thin layers of SiO or KC ℓ sandwiched between layers of Bi (0.6 μm) and mylar ($\sim 10 \mu\text{m}$). On the time scale of the laser pulse this material will not decompress. Furthermore when subjected to a laser pulse on the Bi laser, the tracer layer will be shock compressed. The ablation pressures will be $\sim 10 \text{ Mbar}$ which will produce a quasi strong shock with a compression of $C \sim 3$. The electron densities of the various states of ionisation Z would then be quite well defined at CZn_i where n_i is the ion number density. The electron densities of SiX and C & XIII would then be 1.3×10^{24} and 1.4×10^{24} respectively.

As discussed in 5.3, absorption spectra were taken by viewing the unresolved transition array (UTA) continuum from the Bi plasma through the absorber layer. Typical spectra are shown in Figure 5.06. The wavelength for the Si case was determined from the Si XIII $1^1S - 2^1P$ emission line and the Si K α line. For the case of the Cl absorption spectrum the wavelength scale was fixed by the Cl $1^1S - 2^1P$ emission line and the dispersion of the spectrometer.

Time resolved measurements of these absorption features (5.3) showed that they all appeared at the same time.

Theoretical spectra, from the Cowan code (5.16), are shown in Figure 5.06. Comparison with the experimental spectra shows broadening which is probably due to the high density. A measure of the instrumental width is given by the width of the Cl XVI $1^1S - 2^1P$ emission line seen.

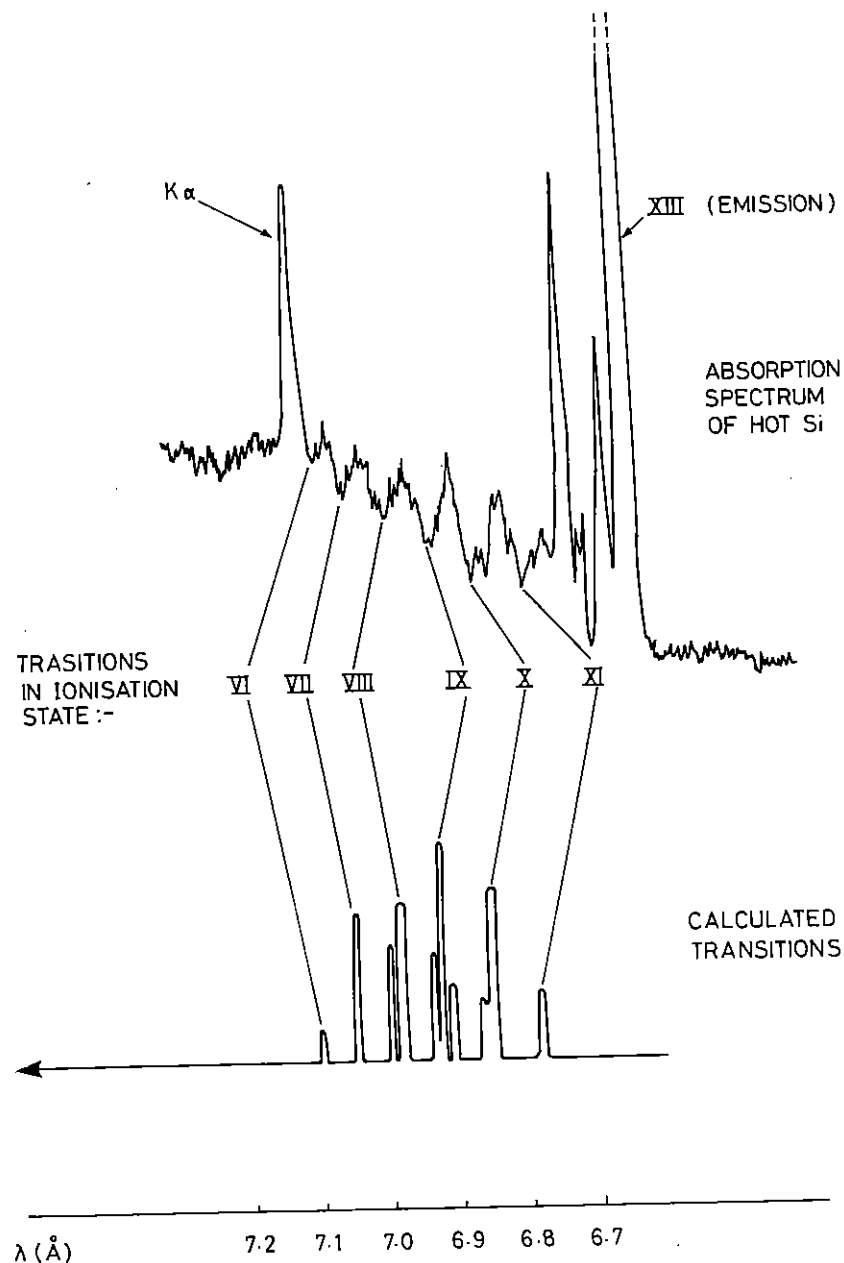
Emission spectra have been obtained at lower density in laser produced plasmas. Figure 5.07 and in tokamaks (5.17) showing similar shifted 1-2 transitions in Ti and Fe. The spectra reported here are simpler in that they are absorption spectra and therefore start from the ground state. A fuller analysis of these data is planned.

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5.5 X-ray Spectroscopy of Low Opacity Polymer Targets

The use of polymer shells, filled with dichloro-difluoromethane in laser compression experiments is described. The absence of shell emission leads to chlorine spectra with little contaminating shell continuum emission. The low X-ray opacity of polymer indicates that chlorine helium-like emission could be detected even with compressed shell areal densities of 0.1 gm cm^{-2} .

Laser compression targets have used glass microballoons with a shell thickness that is practically limited to a minimum of $1 \text{ }\mu\text{m}$, and a fuel



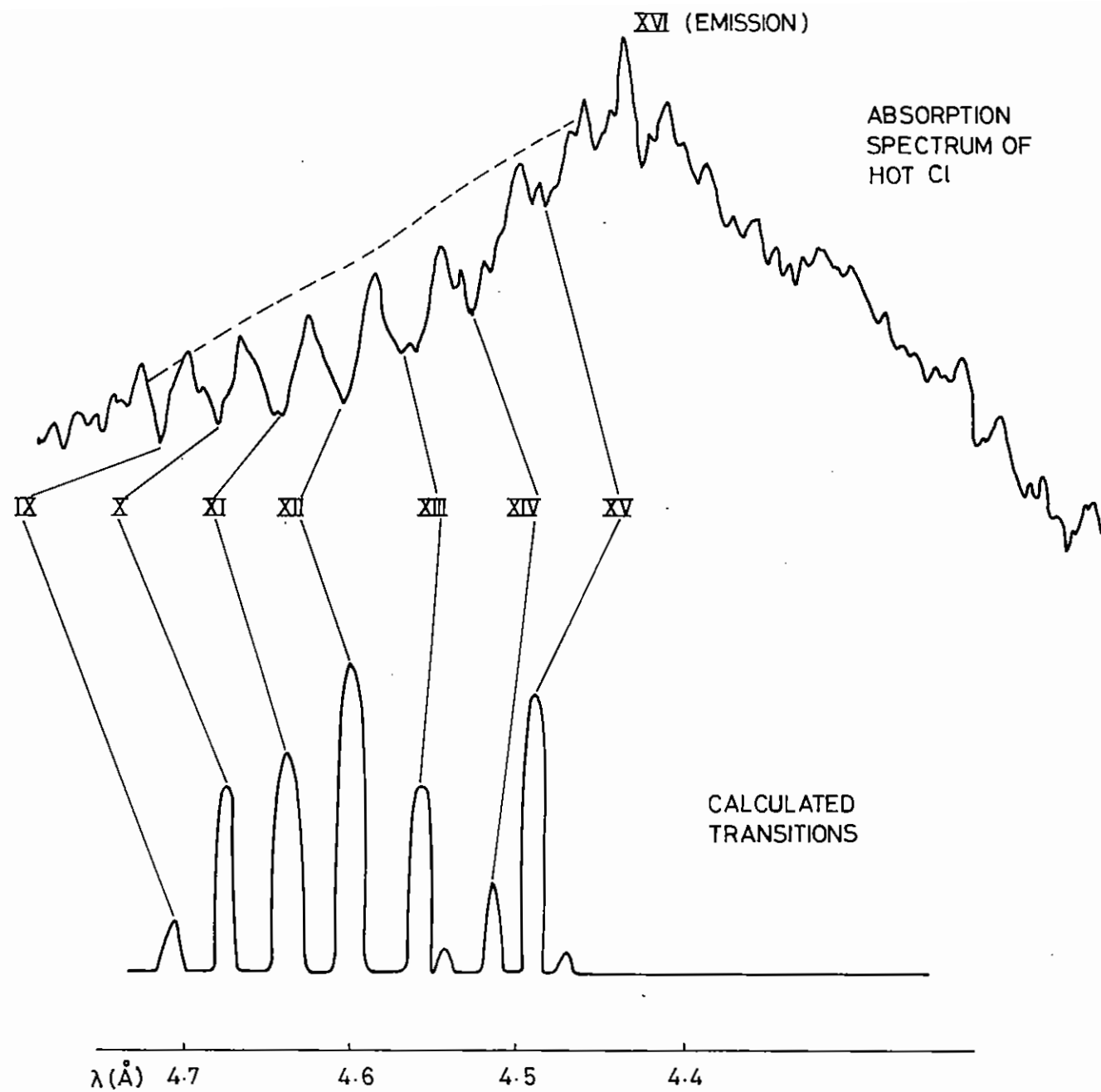


Fig. 5.06 b) Absorption spectra of shock heated Cl.

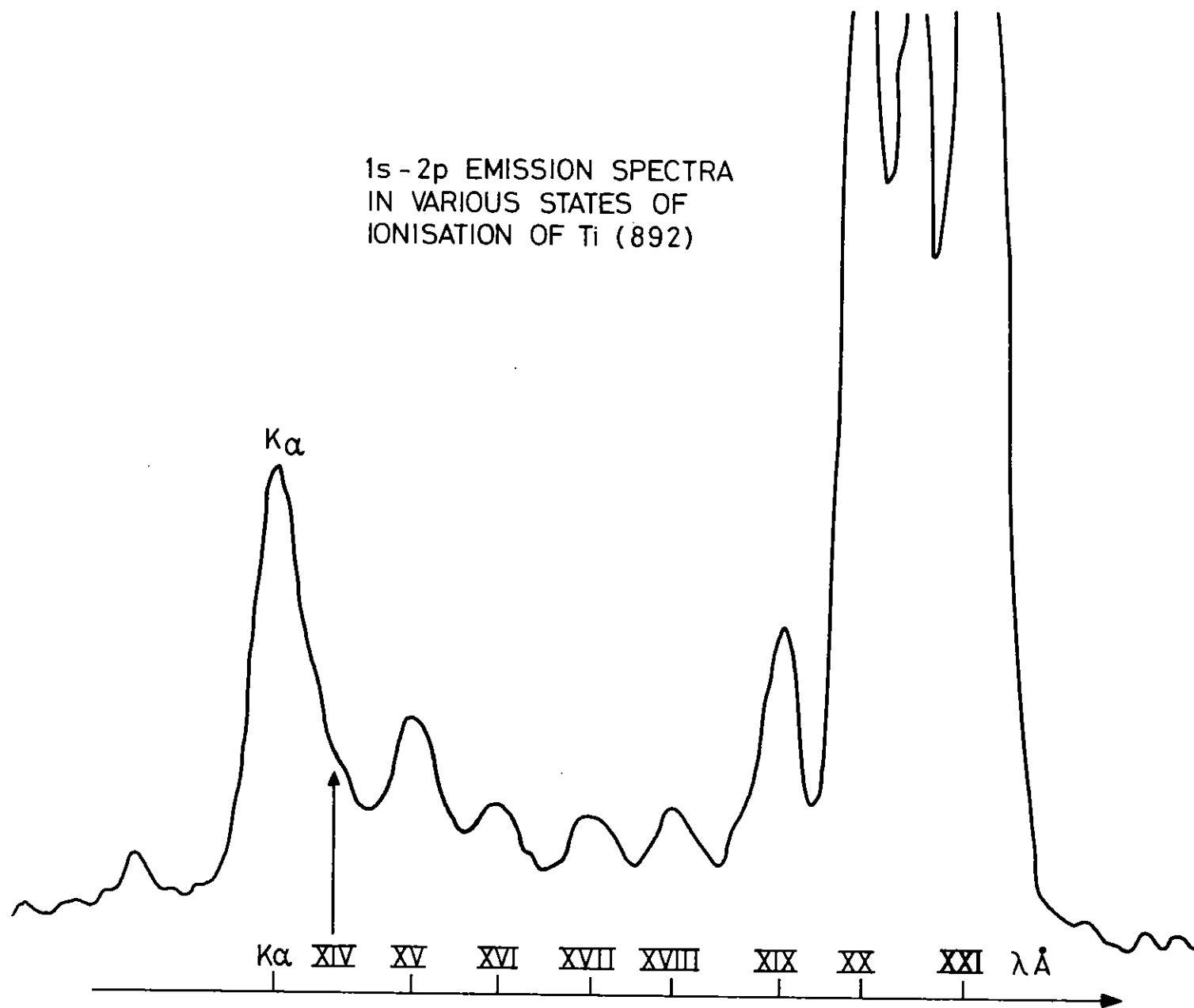


Fig. 5.07 Absorption spectrum of Ti.

doped with a diagnostic gas for X-ray spectroscopy (5.18). In the 'exploding pusher' mode these simple targets readily produce hydrogen and helium-like emission lines from elements up to argon. The photon energy of these lines is high enough that they are not absorbed by the surrounding compressed glass which, in this mode has an areal density of about 10^{-3} gm cm⁻². However in the more desirable ablative compression mode (5.19) higher areal densities of the compressed shell have been achieved (5.08)(5.09)(5.20). Glass shells are opaque to X-rays in the 3 keV region with these targets because of the high X-ray mass attenuation coefficient (μ/ρ) of SiO₂ and the large compressed areal density.

In this paper we describe the manufacture of targets for which the glass shell is replaced by polymer. The cold X-ray mass attenuation coefficient of plastic is a factor of four less than that of glass, and moreover the absorption coefficient of the compressed shell will be further reduced by any temperature rise. Exploding pusher experiments are described with such targets that demonstrate their ease of fill and use. The fill gas is dichlorodifluoromethane (5.21) and diagnosis of the compressed chlorine indicates electron densities typical of exploding pusher experiments. High quality spectra, with little continuum emission below the recombination edges results from the absence of any element of higher atomic number than carbon in the shell. The advantages of these targets for ablative implosions are discussed.

Polymer targets were made by using solvent-drying techniques (5.22). Polystyrene was dissolved in dichloroethane to form a 10% by weight solution. Water was then added to this solution and the mixture beaten to form an emulsion containing water droplets of the desired size. This emulsion was then transferred to another vessel containing water, during which process each water droplet carried over with it an outer layer of polystyrene solution. The mixture was stirred and heated slowly to 50°C to evaporate the dichloroethane (under water) leaving hard water-filled shells. After extraction by sieving the shells were maintained at 40°C for a few days to diffuse out the water. Shells with diameters ranging from 50 to 500 μ m and wall thicknesses from 1 to 20 have been produced by these methods by varying the beating and mixing process. Uniformity and

wall thickness measurements were made using interference microscopy since the shells were transparent. Surface examination using a scanning electron microscope indicated that the shells were very smooth and free of any defects greater than 0.1 μ m in size.

Targets selected to be suitable for implosion experiments were glued to fine (5 μ m diameter) carbon fibres and placed in a vessel which was evacuated and filled with 1 atmosphere of dichlorodifluoromethane (CCl₂F₂).

Measurements of the permeability of polystyrene to various gases were made using films cast from 10% solution in dichloroethane. The results are shown in Table 5.02. These data indicate that for the targets used in these experiments (diameter 120 μ m, wall thickness 2 - 4 μ m) air diffuses from the shells and is replaced by CCl₂F₂ in times of order 2 hours.

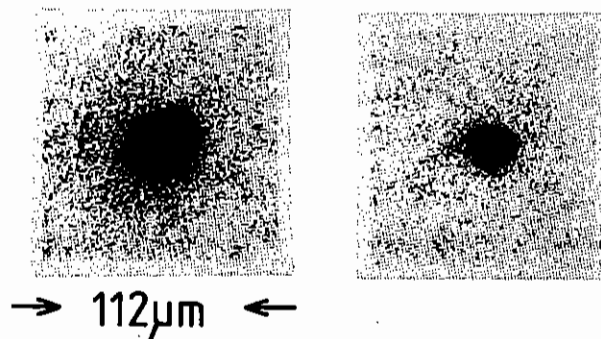
The filled targets were rapidly transferred to the target chamber and aligned in the focus of the six-beam, $\lambda = 1.05$ μ m laser of the Rutherford Appleton Laboratory Central Laser Facility (5.23). Targets were irradiated by 90 ps laser pulses of total energy 80 - 90 J and implosions were thus in the exploding pusher mode where fast electrons dominate the compression (5.18). An X-ray pinhole camera picture of the implosion of a polymer target compared to a similar glass target is shown in Figure 5.08. The absence of any shell emission compared to the glass target is clear.

Time-integrated X-ray spectra were taken with a PET crystal dispersing a spectrum onto Kodak SB2 X-ray film. Coarse microdensitometer scans of the spectra are shown in Figure 5.09 for three target shots with increasing core electron density. Control shots with no fill gas showed no chlorine emission.

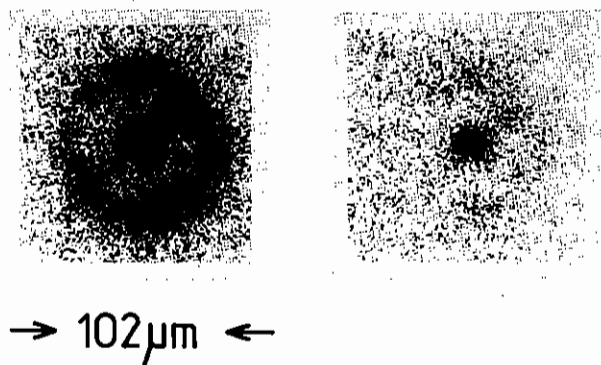
The target parameters are shown in Table 5.03. For comparison, the best argon spectrum from a concurrent set of experiments with argon-filled glass shell is shown in Figure 5.10. There are clear observable advantages of the polymer shell targets viz: (i) the continuum emission at energies below the helium-like recombination edge is considerably less

2 BAR CHLORINE

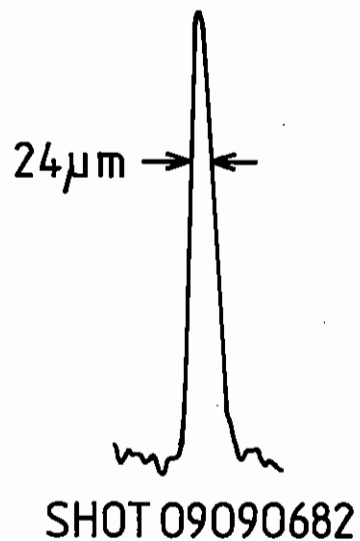
SHOT 18100682



2 BAR ARGON



25 Be + 5 μ m Al 15 μ m Al



Gas	Permeability coefficient (cm ³ (STP). cm/cm ² cm Hg. sec)x 10 ⁻¹⁰
He	2.6 x 10 ⁻¹¹
N ₂	1.0 x 10 ⁻¹²
CCl ₂ F ₂	4.0 x 10 ⁻¹⁴

Table 5.02

Permeability coefficients for various gases through polystyrene films at 20°C.

Fig. 5.08 Pairs of X-ray pinhole camera pictures of plastic targets filled with Freon 12 (upper) and a similar size glass target filled with argon. The filtering on the pinholes is shown.

Shot	Target		E incident J	$\Delta\lambda$ of Cl XVI $1s^2\ ^1S - 1s4p\ ^1P$ mÅ	$n_{e\text{ core}}$ $\times 10^{-23}(\text{cm}^{-3})$	$T_{e\text{ core}}$ eV
	$2r$ μm	Δr μm				
18100682	112	2.3	86	22 ± 2	$1.1 \pm .1$	810
13110682	110	3.7	92	25 ± 2	$1.4 \pm .1$	740
12170682	118	2.9	82	32 ± 3	$2.0 \pm .2$	540

TABLE 5.03 Parameters for the spectra shown in Figure 5.09.

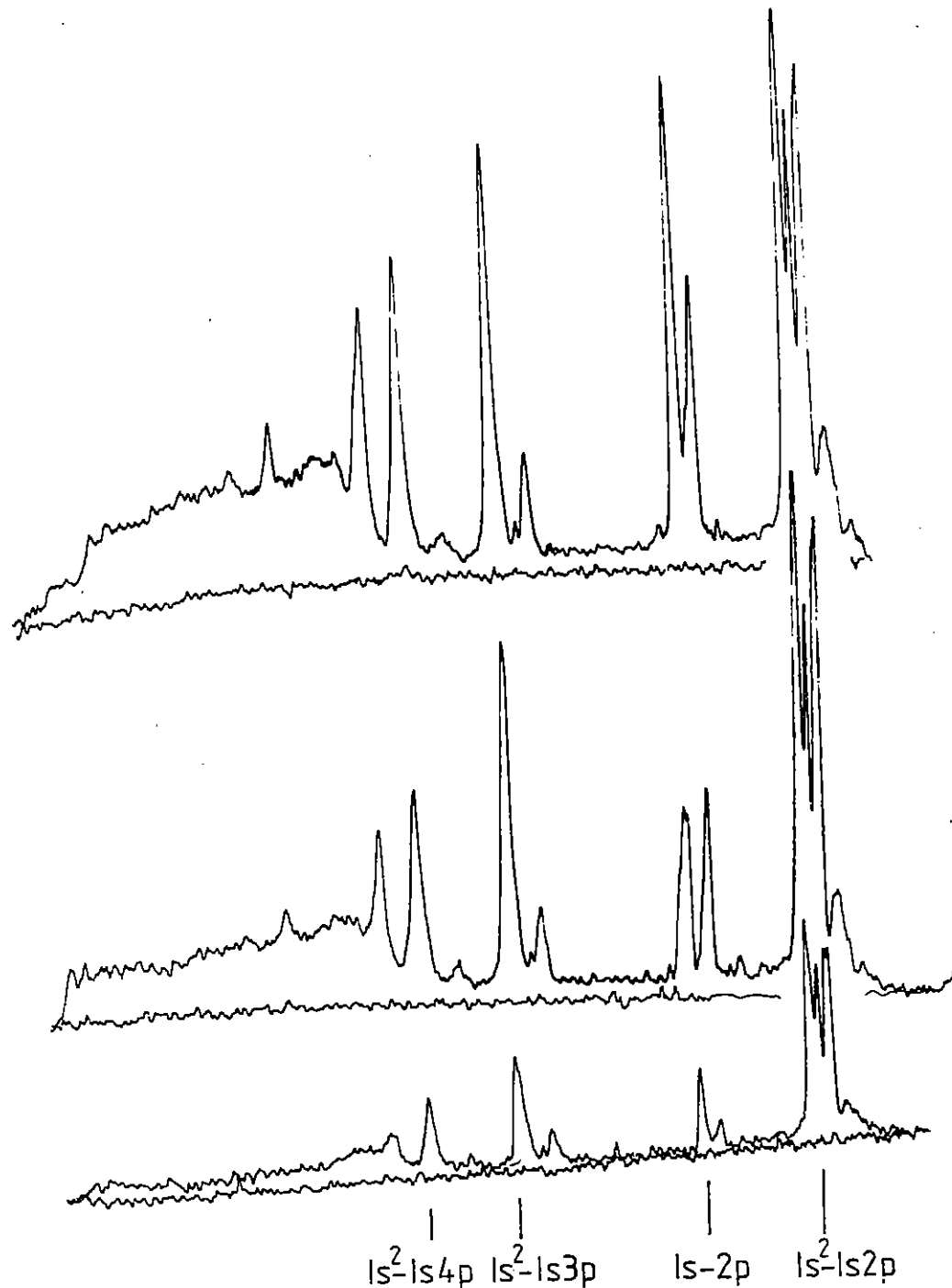


Fig. 5.09 Time integrated spectra of Cl XVI and Cl XVII from plastic targets on shot numbers. 18100682 (top), 13110682 (middle) 12170682 (bottom). The fog level on the film next to the spectrum is also shown.

than for the glass shell, and (ii) that therefore the recombination edge is more pronounced.

Core densities have been obtained from the spectra of Figure 5.09. Expanded traces give the line widths shown in Table 5.03. Electron densities can be obtained from calculation of the Stark broadening (6.24) and are also shown in Table 5.03. The densities can also be obtained from the ratios of the dielectronic satellites of the $1s-2p$ line (Figure 5.11) and the ratios of dielectronic satellites of the $1s^2 \ ^1S - 1s \ 2p^1 \ ^1P$ lines. Additionally, the well-resolved and shifted helium-like recombination edges can be used (5.25).

The temperatures in Table 5.03 are obtained from the ratio of optically thin helium-like to hydrogen-like lines (5.26)(5.27).

The greatest advantage of the polymer shells is in their low X-ray opacity, which is given in Table 5.04 for a variety of temperatures (5.28). For comparison, the opacity of glass (SiO_2) is shown at comparable wavelength. For low temperatures there is a factor of 4 difference, but at higher temperature 100 eV, the opacity of polymer decreases significantly because of ionisation into its K shell.

Present ablative implosions have areal densities of their compressed shells of about $10^{-2} \text{ gm cm}^{-2}$. Glass shells cannot be used for argon emission spectroscopy as the transmission of the argon lines is only 1%. In contrast the transmission of the same area density of polymer at 100 eV is 60% for chlorine lines. If the compressed polymer were further heated to, say, 200 eV by the core emission, significantly higher compressed shell areal densities up to $10^{-1} \text{ gm cm}^{-2}$, could be used with chlorine emission spectroscopy.

Fill gases other than Freon could be used if coatings of lower permeability material were placed on the plastic shell. A layer of low diffusivity material thin enough to entrap the D-T fuel, yet of low enough X-ray opacity, eg $0.1 \text{ } \mu\text{m Al}$, could be coated onto the plastic targets. This could then be overcoated with further plastic to reduce shell

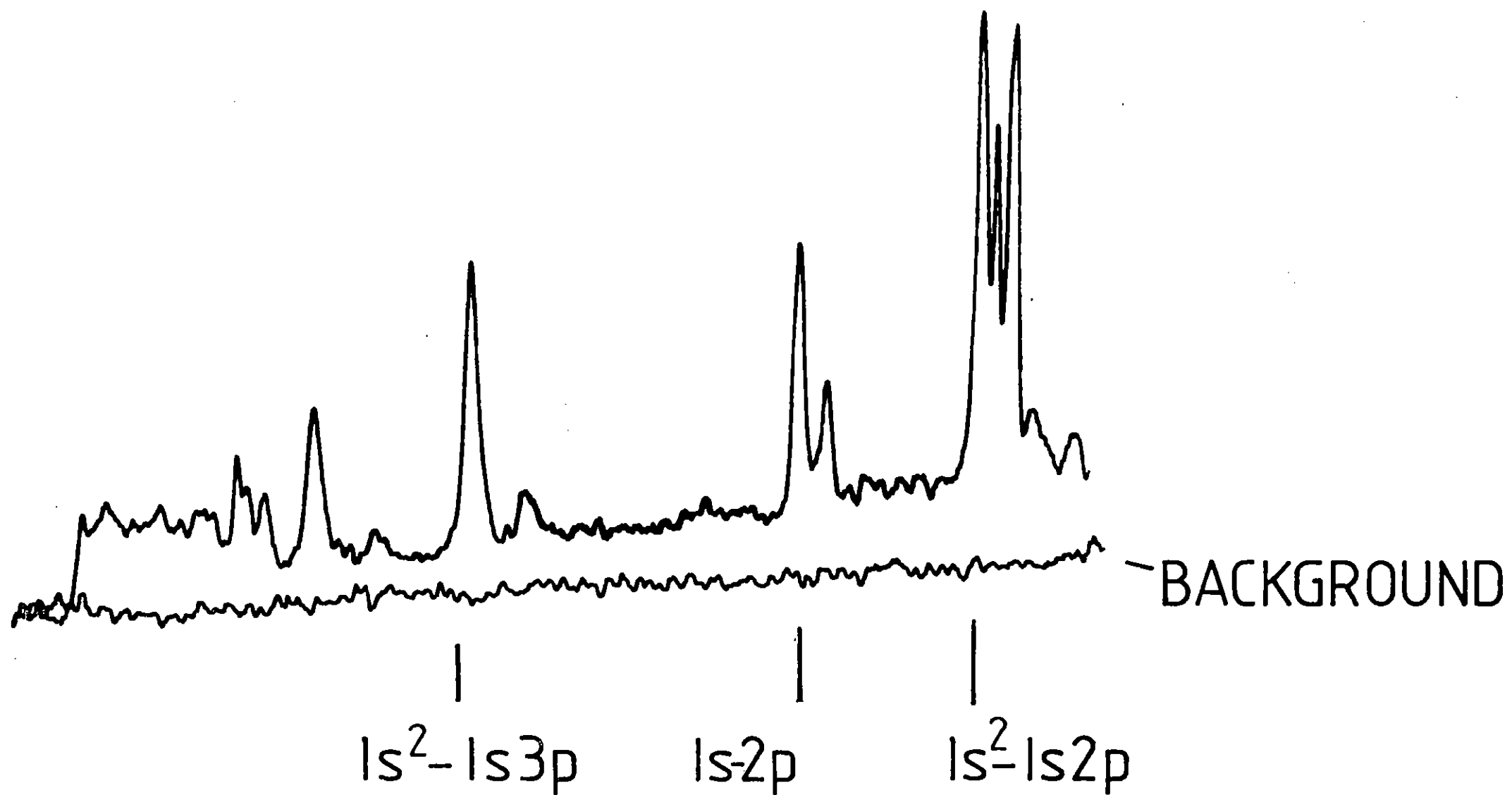


Fig. 5.10 A time integrated spectrum of Ar XVII and Ar XVIII lines for an argon filled glass shell of comparable parameters to the plastic targets in Fig. 2. The Ar XVII recombination edge is visible, although much less pronounced than in Fig. 3.

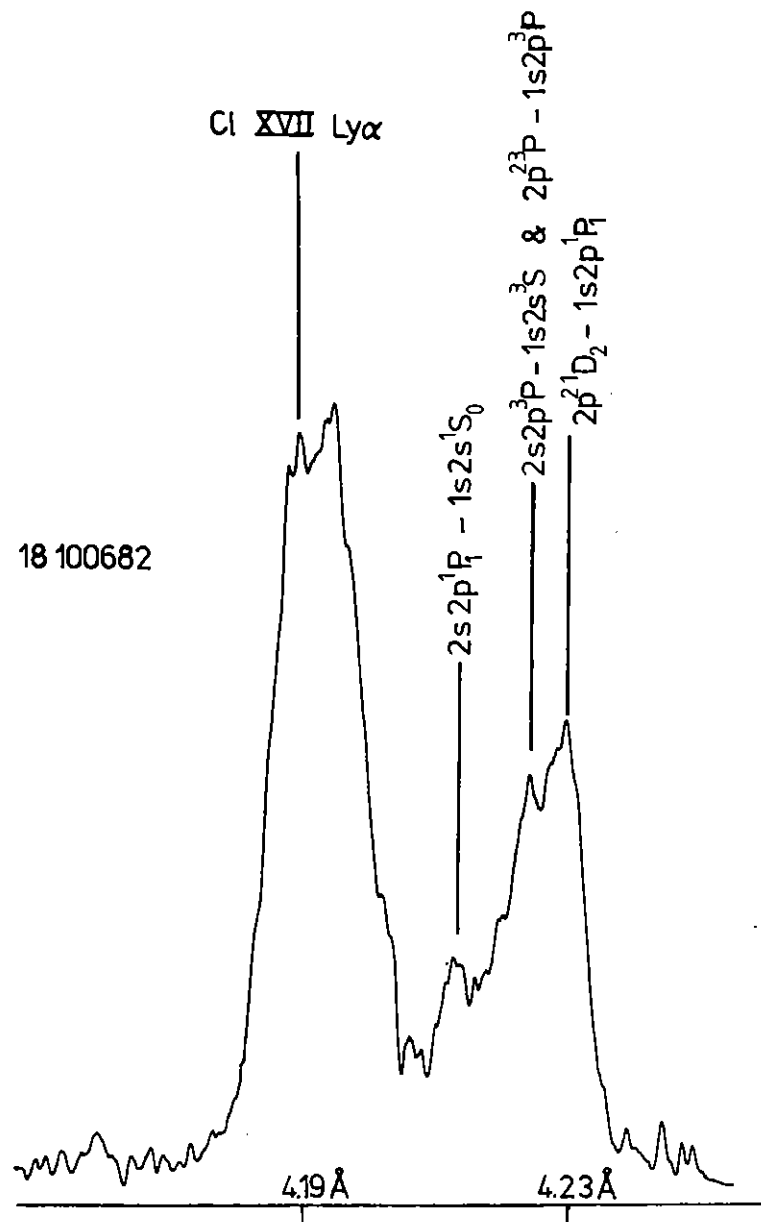


Fig. 5.11 Expanded tracing of the Cl XVII Lyman α transition and its dielectronic satellites from the top tracing of Fig. 1. The satellite triplet-singlet ratio is a density diagnostic and is used in Table 2 to infer electron densities.

Electron temperature T_e (eV)	Mass absorption coefficient (μ/ρ) ($\text{cm}^2 \text{ gm}^{-1}$)			
	Carbon ($h\nu = 2.96 \text{ keV}$)		SiO ₂ ($h\nu = 3.13 \text{ keV}$)	
	$\rho = 1 \text{ gm. cm}^{-3}$	$\rho = 10 \text{ gm. cm}^{-3}$	$\rho = 1 \text{ gm. cm}^{-3}$	$\rho = 10 \text{ gm. cm}^{-3}$
50	113	113	485	488
100	54	89	443	459
200	12	42		
300	6.6	18	337	375
400	4.7	11		

Table 5.04

Variation of the mass absorption coefficient for carbon and glass (SiO₂) with temperature and density. The hydrogen in the polymer will have a negligible effect on the opacity of polymer shells.

emission.

In conclusion we have described a new target fabrication technique that demonstrably works and could be used for ablative implosions.

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5.6 Thermal Smoothing Effects under Steady-State Conditions

5.6.1 Experimental Investigation of Long Pulse Smoothing

Previous experiments using short pulse X-ray backlighting to observe the target smoothing of an imposed laser beam modulation (5.29)(5.30), indicated a scaling with laser irradiance and wavelength of the observed smoothing like $(I\lambda_L)^{1.5}$. Numerical simulations showed that because the measurements were made after the end of the laser pulse, target decompression significantly affected the results.

In order to improve on these experiments and attempt to make an unambiguous measurement of the effects of thermal smoothing further experiments were performed. Making use of the synchronised long/short pulse oscillators developed in the interim period (5.31) it was possible to drive a microdisc target with a smooth 1 ns long laser pulse. A second synchronised 100 ps pulse enabled short pulse X-ray shadowgraphy to be carried out, as before.

A major improvement in the experiment is thus the increased ratio of $\tau_L:\tau_{BL}$, of about 10:1. The long drive pulse is better characterised in its temporal power profile than the stacked pulse previously used. Also the longer irradiation time allows an almost steady-state situation to be reached in terms of plasma flow velocity and scalelength such that the separation of the critical and ablation surfaces is determined by the hydrodynamics rather than by the laser pulse duration, over a significant time interval.

A fixed backlighting delay time was employed, the delay being set so that

the target was shadowgraphed at approximately the end of the drive laser pulse. This ensured that the observation was made during the steady-state phase and before the end of laser irradiation, whilst the target was still being driven.

A consequence of the slower rising edge of the long drive laser pulse compared to the previously used stacked pulse (particularly for $\lambda_L = 1.05 \mu\text{m}$) is that the initial pressure exerted on the targets is lower than before. To ensure that the targets covered a sufficiently large distance by the observation time to allow an accurate measurement of their position it was found to be necessary to reduce the target thickness by 1.05 μm irradiation to 3.5 μm of Mylar. The higher absorption and increased ablation pressure associated with 0.53 μm irradiation allowed a target thickness of 12.5 μm of Mylar to be used, as previously.

The line of sight mass density through the accelerated region of the thin 3.5 μm targets is reduced from that of the thicker 12.5 μm targets. To ensure that the accelerated target material did not become transparent to the backlighting X-rays filtering was reduced to 12.5 μm of Mylar in the backlighting target substrate with 15 μm of Beryllium over the recording pinhole camera. This resulted in the 1/e cut-off energy of the backlighting X-rays being softened to 1.8 keV. Also the rear surfaces of the targets were coated with a 50 nm layer of Bismuth (equivalent to $\approx$ 0.5 μm of Mylar by mass) to increase the attenuation of the rear side and give a sharper edge to the shadowgrams.

The basic experimental arrangement was the same as that used in the earlier experiments and is shown in Figure 5.12. Shadowgrams were obtained for two laser wavelengths as before, the drive pulse lengths being 1.2 ns and 0.9 ns FWHM for $\lambda_L = 1.05 \mu\text{m}$ and 0.53 μm , respectively. Using an X-ray streak camera with similar filtering the backlighting X-ray pulse length was measured as 100 ps FWHM.

The backlighting pulse was timed to be coincident with the peak of the 1.05 μm main drive pulse and then delayed by 0.5 ns. This resulted in an observation time of 1.3 ns after the start of the drive pulse for

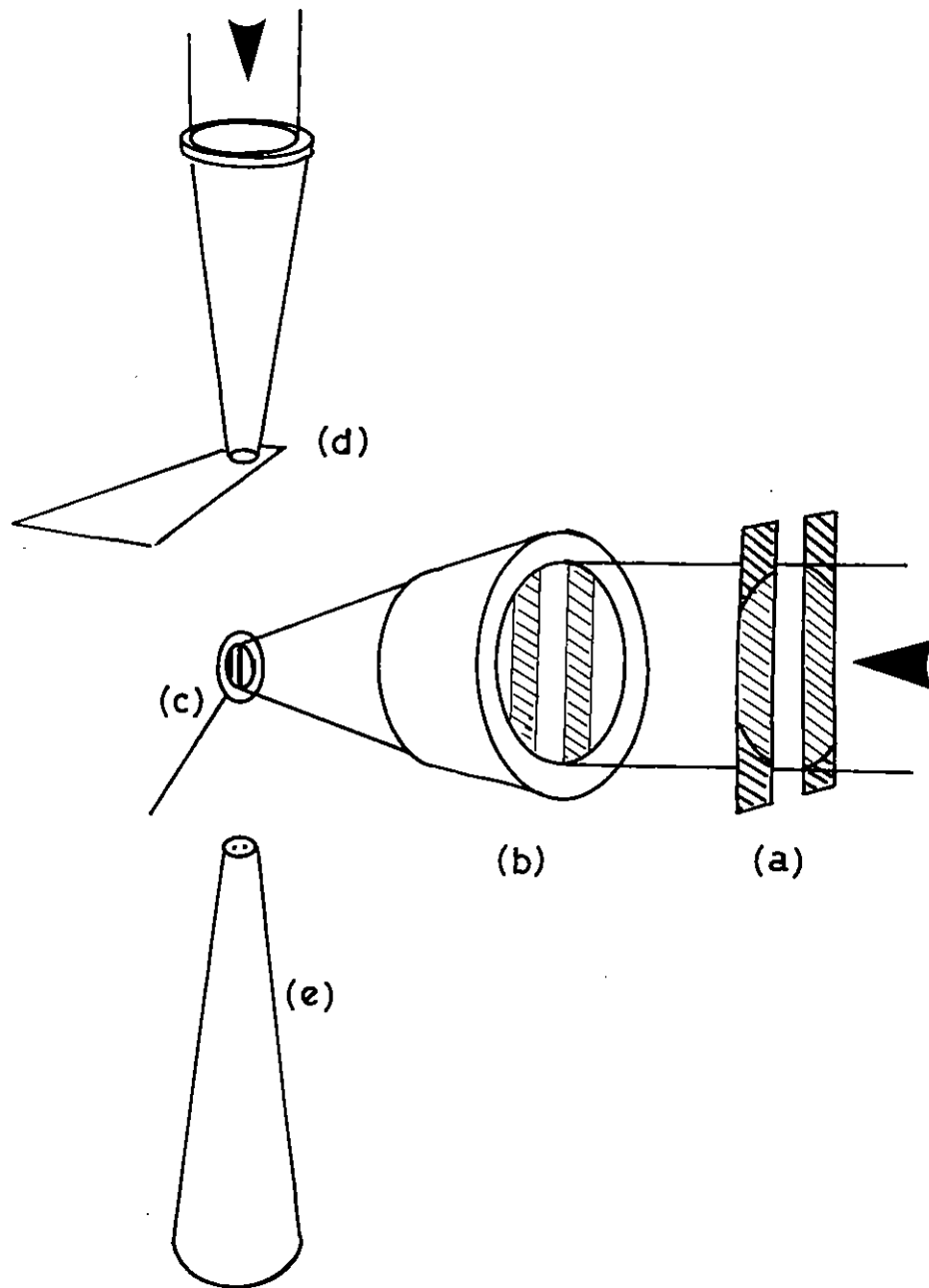


Fig. 5.12 Experimental arrangement showing a) beam modulation, b) $f/1$ focusing lens, c) microdisc target, d) backlighting target and e) recording pinhole camera.

$\lambda_L = 1.05 \mu\text{m}$. The observation time for green irradiation was set at 1.5 ns after the start of the pulse. Shadowgrams were obtained on Kodak SB2 X-ray film using a pinhole camera equipped with a $10 \mu\text{m}$ platinum pinhole.

The irradiance modulation depth was reduced to 2:1 to be similar to the irradiance variation on a spherical target in a multibeam compression experiment. A fixed focusing condition was used of $250 \mu\text{m}$ beyond the target surface, in nearly all the experiments. This resulted in a modulation wavelength λ_p of $90 \mu\text{m}$ in the target plane. To achieve the highest green irradiances it was necessary to decrease the nominal focal spot size to $200 \mu\text{m}$, reducing λ_p to $75 \mu\text{m}$ in the target plane.

5.6.2 Results

Four shadowgrams obtained with $1.05 \mu\text{m}$ irradiation are shown in Figure 5.13. The rear surface distance modulation can be seen to decrease as the incident irradiance is increased, in a similar way to that seen in the earlier experiment. As outlined above, all the results were obtained with a fixed delay time so the distance travelled also increases with increasing irradiance, as Figure 5.13 shows.

The shadowgrams for both laser wavelengths were measured by microdensitometry across the high and low drive regions and the distance modulation Δ (defined as $\frac{\text{Maximum distance travelled}}{\text{Minimum distance travelled}} - 1$), corrected for initial target thickness, and peak and trough pressures were determined. The data are listed in Table 5.05.

The distance modulation results are shown in Figure (5.14) as a function of the absorbed laser irradiance, calculated using the absorption data given by Evans et al (5.32). As before the measurement accuracy increases as the irradiance is increased and the distance travelled by the target gets larger. At low irradiance the target modulation reflects the imposed irradiance modulation. The value of Δ at low irradiance is 0.5–0.6 which agrees well with a pressure-irradiance relationship like $P \propto I^{0.7}$ and a

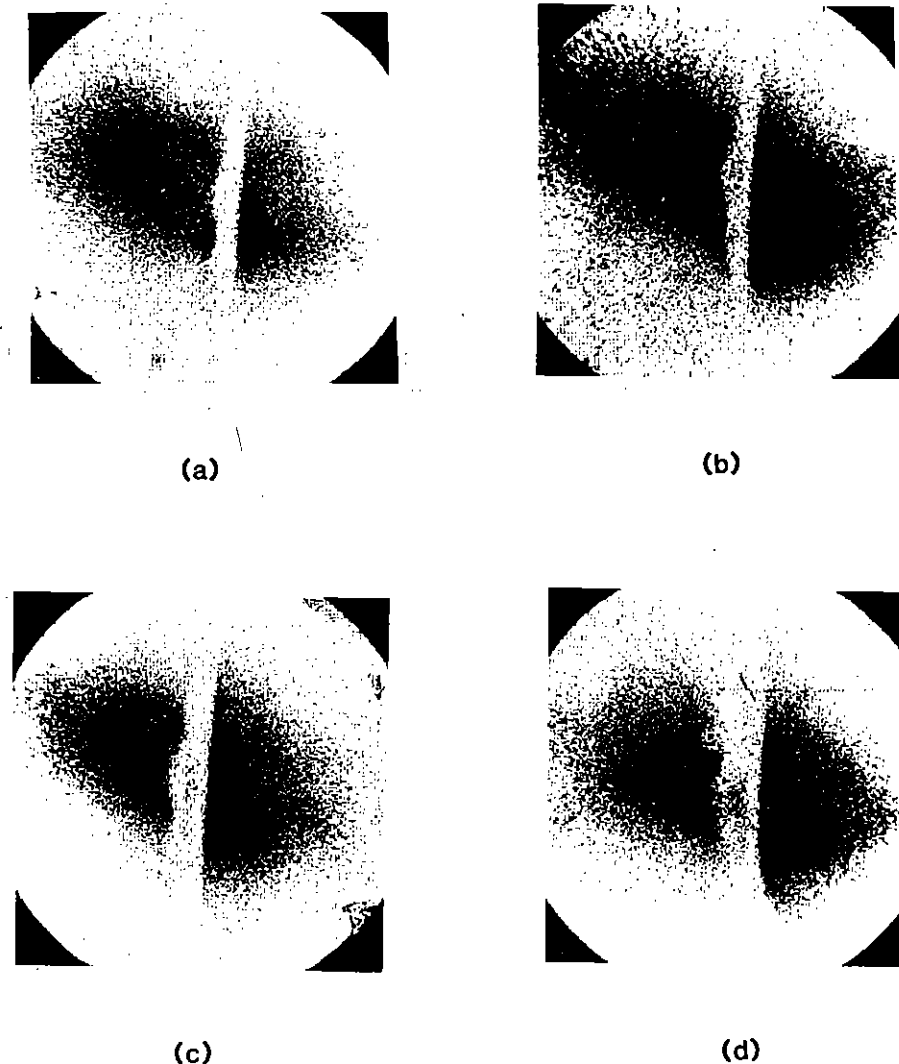


Fig. 5.13 Shadowgrams for $\lambda_L = 1.05 \mu\text{m}$, long pulse, and a range of absorbed irradiances.

- a) $I_{\text{abs}} = 7.4 \times 10^{12} \text{Wcm}^{-2}$, $250 \mu\text{m}$ focal spot, $\lambda_p = 90 \mu\text{m}$, $\Delta = 0.50$.
- b) $I_{\text{abs}} = 1.0 \times 10^{13} \text{Wcm}^{-2}$, $250 \mu\text{m}$ focal spot, $\lambda_p = 90 \mu\text{m}$, $\Delta = 0.56$.
- c) $I_{\text{abs}} = 2.1 \times 10^{13} \text{Wcm}^{-2}$, $250 \mu\text{m}$ focal spot, $\lambda_p = 90 \mu\text{m}$, $\Delta = 0.44$.
- d) $I_{\text{abs}} = 3.1 \times 10^{13} \text{Wcm}^{-2}$, $250 \mu\text{m}$ focal spot, $\lambda_p = 90 \mu\text{m}$, $\Delta = 0.10$.

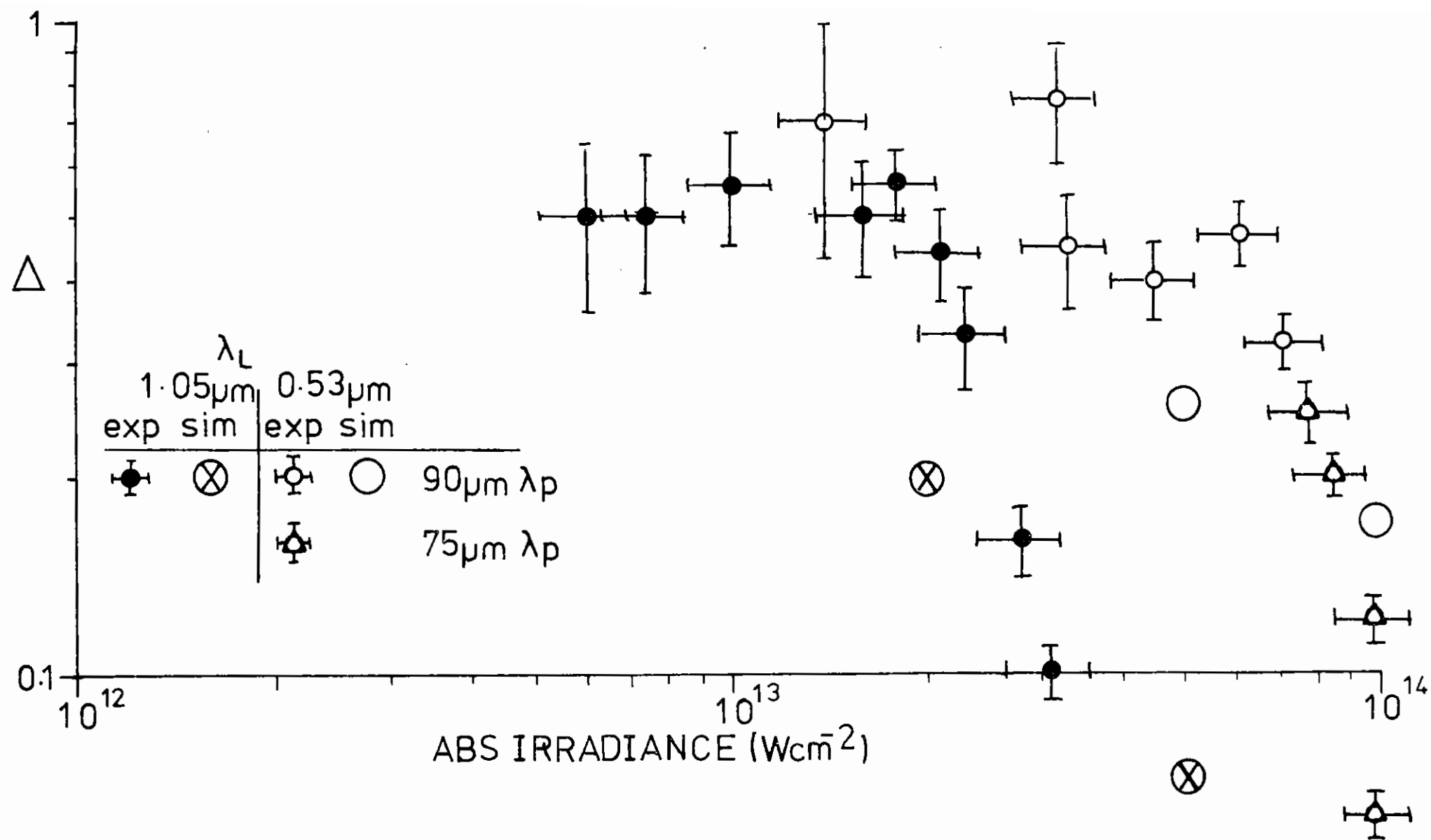


Fig. 5.14 $\Delta = (Z_{\text{MAX}}/Z_{\text{MIN}}) - 1$ versus absorbed irradiance.

Table 5.05

$$\lambda_L = 1.05 \mu\text{m}$$

Shot	Energy J	I_{abs} $\times 10^{13} \text{ W cm}^{-2}$	Focus μm	λ_p μm	ZMAX μm	ZMIN μm	$\frac{Z_{\text{MAX}}}{Z_{\text{MIN}}} - 1$	PMAX Mb	PMIN Mb
29/6:13	13.1	1.6	250	90	43.4	29.0	0.50	3.8	2.6
30/6:06	24.6	2.8	250	90	52.2	45.0	0.16	4.6	4.0
30/6:08	7.7	1.0	250	90	40.6	26.1	0.56	3.6	2.3
1/7:02	5.4	0.74	250	90	34.8	23.2	0.50	3.1	2.0
1/7:03	4.2	0.60	250	90	30.4	20.3	0.50	2.7	1.8
1/7:04	16.9	2.1	250	90	52.2	36.2	0.44	4.6	3.2
1/7:05	20.0	2.3	250	90	46.4	34.8	0.33	4.1	3.1
1/7:07	28.5	3.1	250	90	81.2	75.4	0.10	7.1	6.0
1/7:08	15.4	1.8	250	90	63.8	40.6	0.57	5.6	3.6

$$\lambda_L = 0.53 \mu\text{m}$$

26/7:01	~1.4	250	90	21.8	13.1	0.70	4.4	2.7	
26/7:04	11.6	3.3	250	90	42.1	29.0	0.45	8.5	5.8
28/7:12	16.2	4.5	250	90	58.0	40.6	0.40	11.6	8.2
28/7:18	22.2	6.1	250	90	72.5	49.3	0.47	14.2	9.9
28/7:21	10.8	3.2	250	90	43.5	24.7	0.76	8.7	5.0
29/7:03	24.1	9.7	200	75	81.2	76.9	0.06	16.0	15.5
29/7:08	18.9	7.8	200	75	78.3	62.4	0.25	15.7	12.5
29/7:10	24.5	9.9	200	75	89.9	79.8	0.12	17.8	16.0
29/7:16	20.8	8.5	200	75	94.2	78.3	0.20	18.8	15.6
29/7:19	25.4	10.3	200	75	94.2	92.8	0.02	18.8	18.6
29/7:20	26.6	7.1	250	90	84.1	63.8	0.32	16.8	12.8

2:1 irradiance modulation.

Above an irradiance of $2 \times 10^{13} \text{ W cm}^{-2}$ the $1.05 \mu\text{m}$ results start to show a significant decrease in the degree of target modulation. For $\lambda_L = 0.53 \mu\text{m}$ this onset of smoothing occurs above $I_a = 6 \times 10^{13} \text{ W cm}^{-2}$. The small change in λ_p for the high irradiance green data does not appear to have had a large effect. Possibly the decrease in focal spot size caused an accompanying decrease in the plasma scalelengths, thus keeping the value of D/λ_p constant.

Analytic fits to the data above the threshold irradiances for smoothing yield the following approximate, empirical, relationship:

$$1/\Delta = 0.37(I_a \lambda_L^2)^3$$

I_a in units of $10^{13} \text{ W cm}^{-2}$
 λ_L in units of $1 \mu\text{m}$

This scaling of smoothing with laser wavelength is in good agreement with numerical simulations (5.33) where a λ_L^2 relationship was shown for both the scaling D and the instantaneous pressure modulation. However the high scaling with absorbed irradiance does not match with any numerical studies.

The pressure data is shown in Figures 5.15 and 5.16 for the two laser wavelengths. An allowance has been made for mass loss by laser ablation which for the thin targets is about 25% of the initial mass. A comparison has been made with the data of Goldsack et al (5.14) and good agreement can be seen between this data and the peak pressure points at low irradiance for both laser wavelengths. For the $1.05 \mu\text{m}$ data (Figure 5.15) the pressure drops below the unperturbed level at about $2 \times 10^{13} \text{ W cm}^{-2}$ and for the $0.53 \mu\text{m}$ data (Figure 5.16) at about $8 \times 10^{13} \text{ W cm}^{-2}$. However, particularly for the green data, the fall-off is less than might be expected from the degree of smoothing indicated in Figure 5.14 for $I_a > 8 \times 10^{13} \text{ W cm}^{-2}$. This implies that the total distance travelled by the target is too great and may indicate some form of decompression effect, possibly caused by preheating at these high

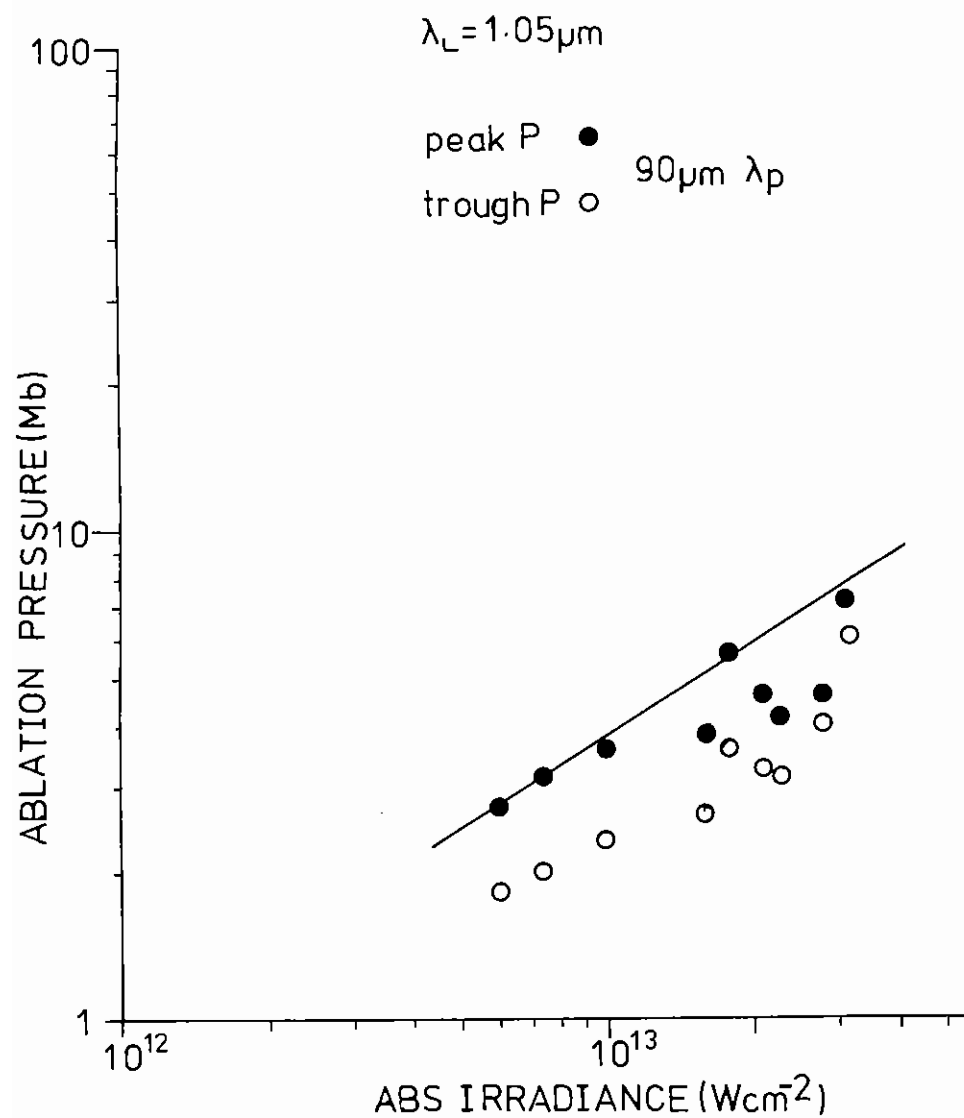


Fig. 5.15 Pressure versus absorbed irradiance for $\lambda_L = 1.05 \mu\text{m}$.

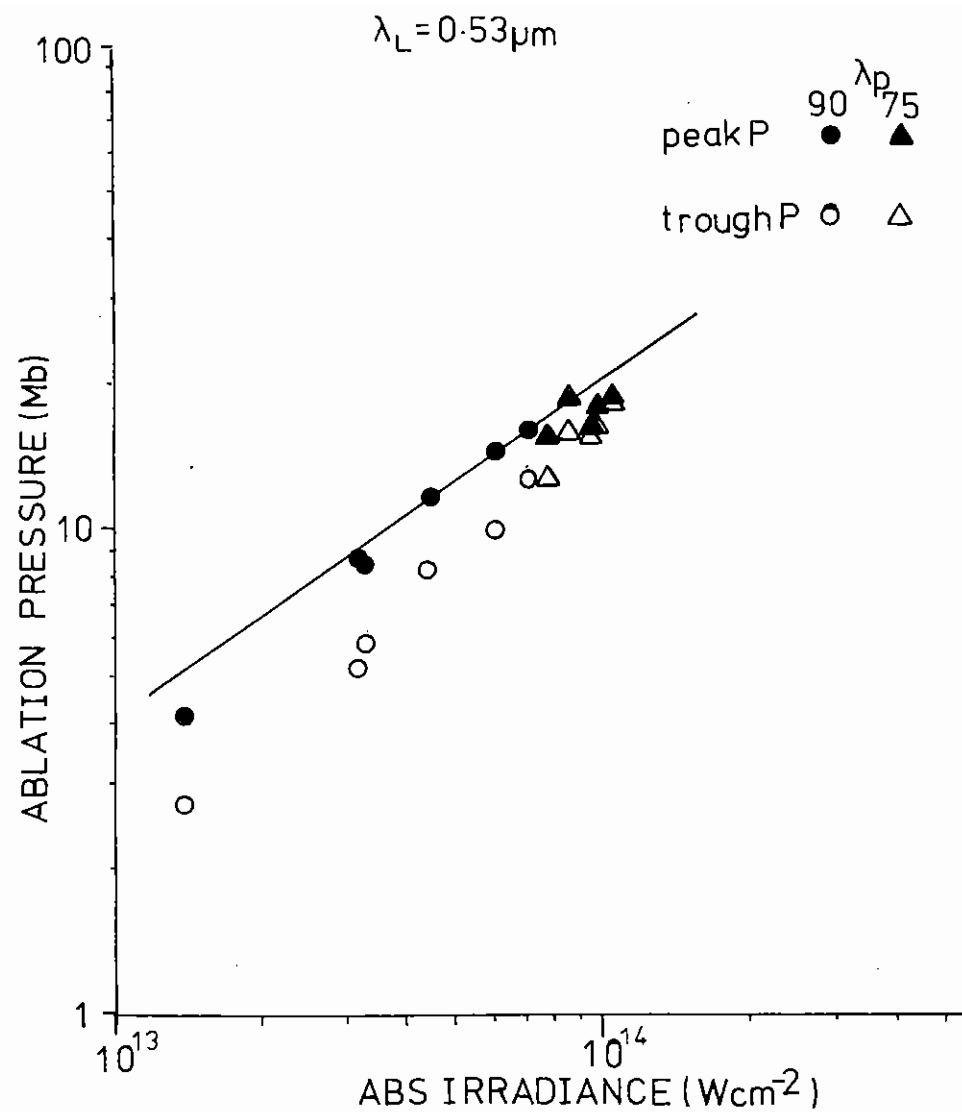


Fig. 5.16 Pressure versus absorbed irradiance for $\lambda_L = 0.53 \mu\text{m}$.

irradiances. This may also explain the strong scaling of the observed smoothing with absorbed laser irradiance with some of the smoothing being a result of decompression in the high density region of the target.

5.6.3 Simulations of Long Pulse Smoothing

The 2-dimensional fluid code Pollux (5.34) was modified to simulate the long pulse experiment described above. Two simulations were performed with $\lambda_L = 1.05 \mu\text{m}$ and a $4 \mu\text{m}$ thick plastic target at absorbed irradiances of 2×10^{13} and $5 \times 10^{13} \text{ W cm}^{-2}$. For $\lambda_L = 0.53 \mu\text{m}$ a target thickness of $12 \mu\text{m}$ was used and simulations performed at absorbed irradiances of 5×10^{13} and $1 \times 10^{14} \text{ W cm}^{-2}$.

In Figure 5.17 the trajectories of the critical and ablation surfaces for the four simulations are shown. The scaling of D can be determined as:

$$D \propto (I_a \lambda_L^3)^{0.9}$$

This is mid-way between the short pulse result of $(I_a \lambda_L^2)^{0.6}$ (5.33) and the 1-dimensional simulation result of Gardner and Bodner (5.35) of

$$(I_a \lambda_L^{3.8})^{0.7}.$$

A plot of the distance modulation as a function of time for the $1.05 \mu\text{m}$ simulations is shown in Figure 5.18. At late time ($> 1 \text{ ns}$) the value of Δ is nearly constant indicating a nearly steady-state situation has been established. As can be seen, the higher irradiance case shows less remaining target modulation than the low irradiance case.

The degree of target modulation at the experimental observation time is shown in Figure 5.14. These simulation results straddle the experimental data and agree with the experimentally determined wavelength scaling for smoothing of λ_L^2 . However a much slower variation with absorbed irradiance is indicated by the simulations, the overall scaling being:

$$1/\Delta = 2.7(I_a \lambda_L^2)^{0.9}$$

5.19

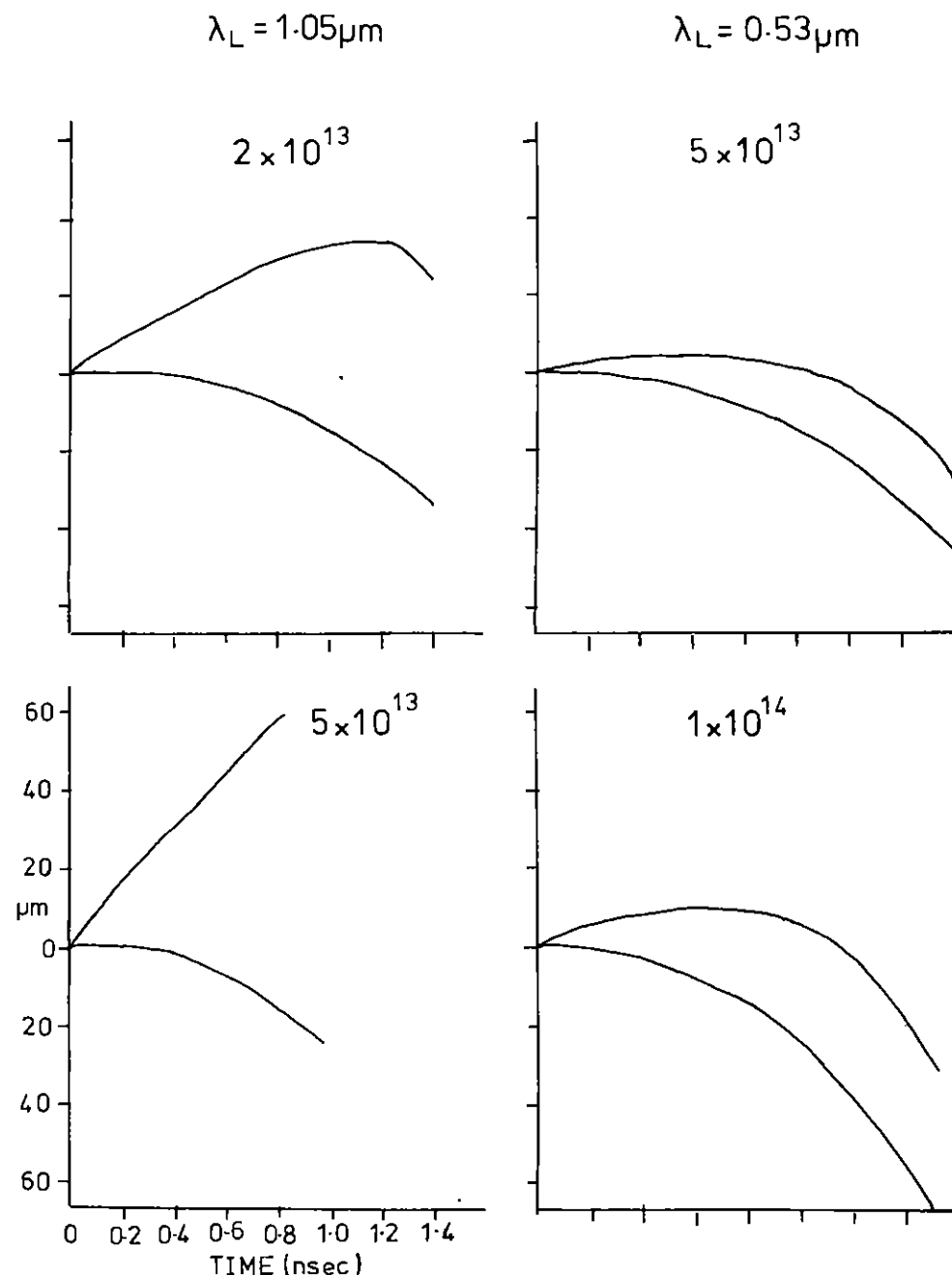


Fig. 5.17 Simulations of critical and ablation trajectories.

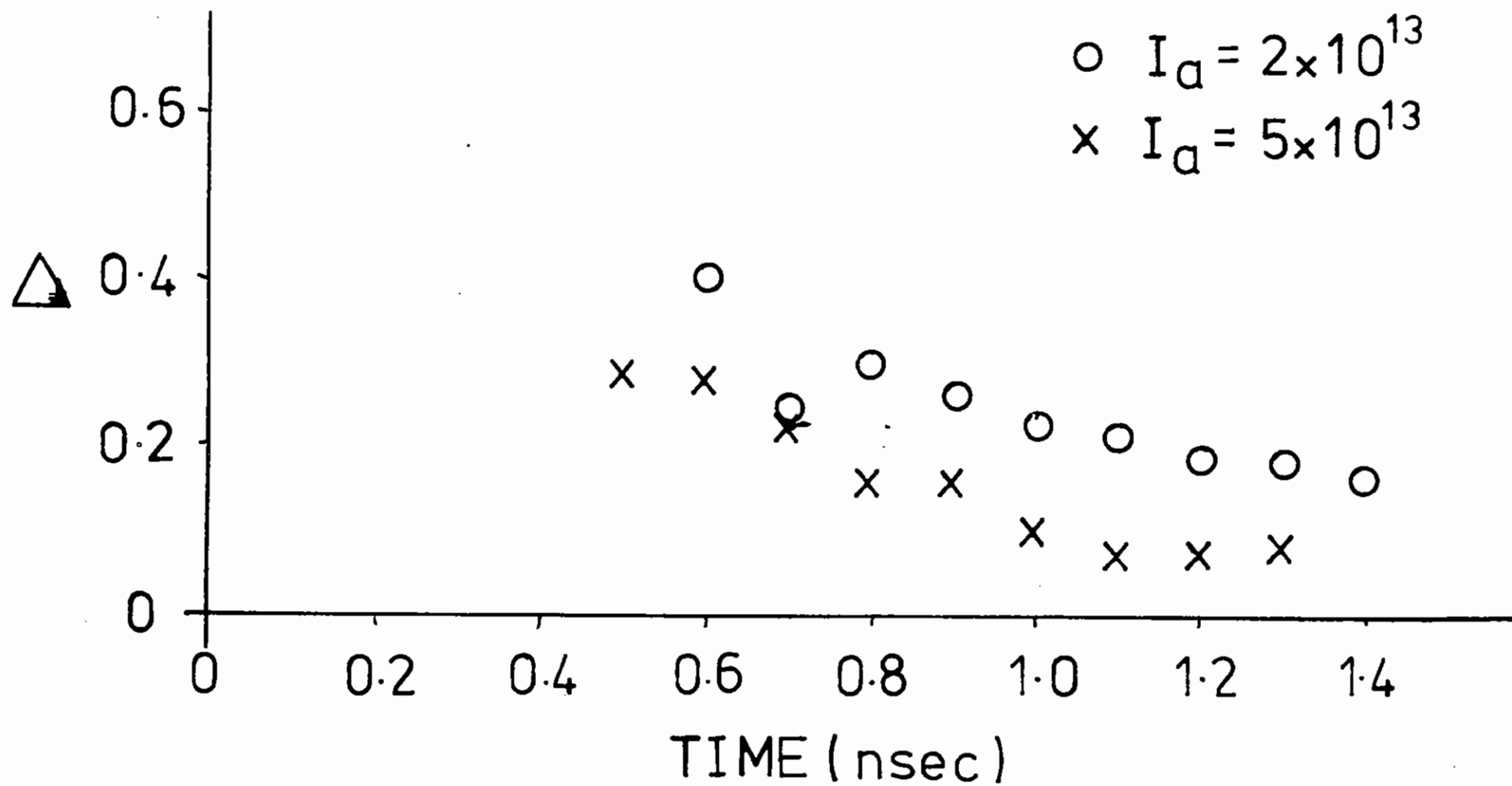


Fig. 5.18 Simulations of $\Delta = (Z_{\text{MAX}}/Z_{\text{MIN}}) - 1$ versus time for $\lambda_L = 1.05 \mu\text{m}$.

$$\begin{array}{l} I_a \text{ in units of } 10^{13} \text{ W cm}^{-2} \\ \lambda_L \text{ in units of } 1 \mu\text{m} \end{array}$$

The irradiance scaling matches that found for the critical to ablation separation distance D given above. The target smoothing however appears to scale more slowly with laser wavelength than does the stand-off distance D .

The simulation results also indicate that at high irradiance ($> 5 \times 10^{13} \text{ W cm}^{-2}$) the measured distances travelled by the targets in the experiments were too great by up to 15%. It therefore seems likely that the strongly driven targets were experiencing some decompression and expansion.

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5.7 Thermal Smoothing Measurements by Shock Transit Time

In the past two years considerable experimental and theoretical evidence (5.35-5.37) has been presented on the use of a heat diffusive mechanism to improve implosion symmetry. In this process the inherent laser beam inhomogeneities are 'thermally' smoothed out by lateral diffusion of energy driving transport from the critical density surface to the ablation front thus tending to improve the effective beam quality by spatially integrating the variations in laser beam profile. Of recent interest is how such effects might scale with change of laser wavelength.

Previous studies of the effects of imposed laser beam non uniformities on ablatively accelerated targets have been performed using two distinct techniques: the direct lateral observation of target modulation with X-ray shadowgraphy using a backlighting beam (5.29)(5.30) at both 0.53 and 1.06 μm and optical streak photography of non uniform blackbody emission from the rear surface of double foil targets (5.38) at 1.06 μm . However in both these studies practical difficulties have prevented measurements on a sub nanosecond timescale. Computer simulations (5.29)(5.30) have shown

that the degree of smoothing observed may be anomalously high due to effects such as hydrodynamic flow and target decompression during a coasting phase.

The experiment described here uses a different approach to measure the smoothing of imposed beam modulations early in time (typically, at 0.5 ns). Moreover, the wavelength of irradiation used is $\lambda_L = 0.35 \mu\text{m}$ where previously no experimental data has yet been published.

It is well understood that during target irradiation material ablates off drilling a shockwave forward into the target, heating the foil material as it goes. As irradiation continues the pressure rises and subsequent shockwaves overtake the preceeding ones to form a propagating shock front. As the shockfront approaches to within a few skin depths of the rear surface of a plane target, the intensity of blackbody emission from the target rises rapidly, reaches a maximum, as the shockfront breaks out and thereafter falls slowly. The spatial profile of this emission corresponds directly to any spatial variations in peak drive pressure on the ablation front. Hence it is possible to measure a parameter which describes the degree of modulation on a shockfront profile introduced by imposed laser beam inhomogeneities of known spatial scalelength and intensity modulations. It is important, however that in order to reliably interpret the degree of modulation observed with an imposed beam non uniformly, the profile of the shockfront breakout from an unmodulated beam should be reasonably planar and have no significant structure on the shockfront profile. In the experiment this requirement has been found to be adequately fulfilled.

Experiment

The experimental layout is shown in Figure 5.19. A single beam of the VULCAN Nd glass laser system (5.23) was frequency tripled to give up to 30 J in 750 ps FWHM. A small fraction of the beam was brought through fibre optics to provide a timing fiducial on the streak photographs. The intensity of the laser beam was modulated by placing 2 cm wide strips of neutral density filter in the 10 cm wide incident beam. The beam was

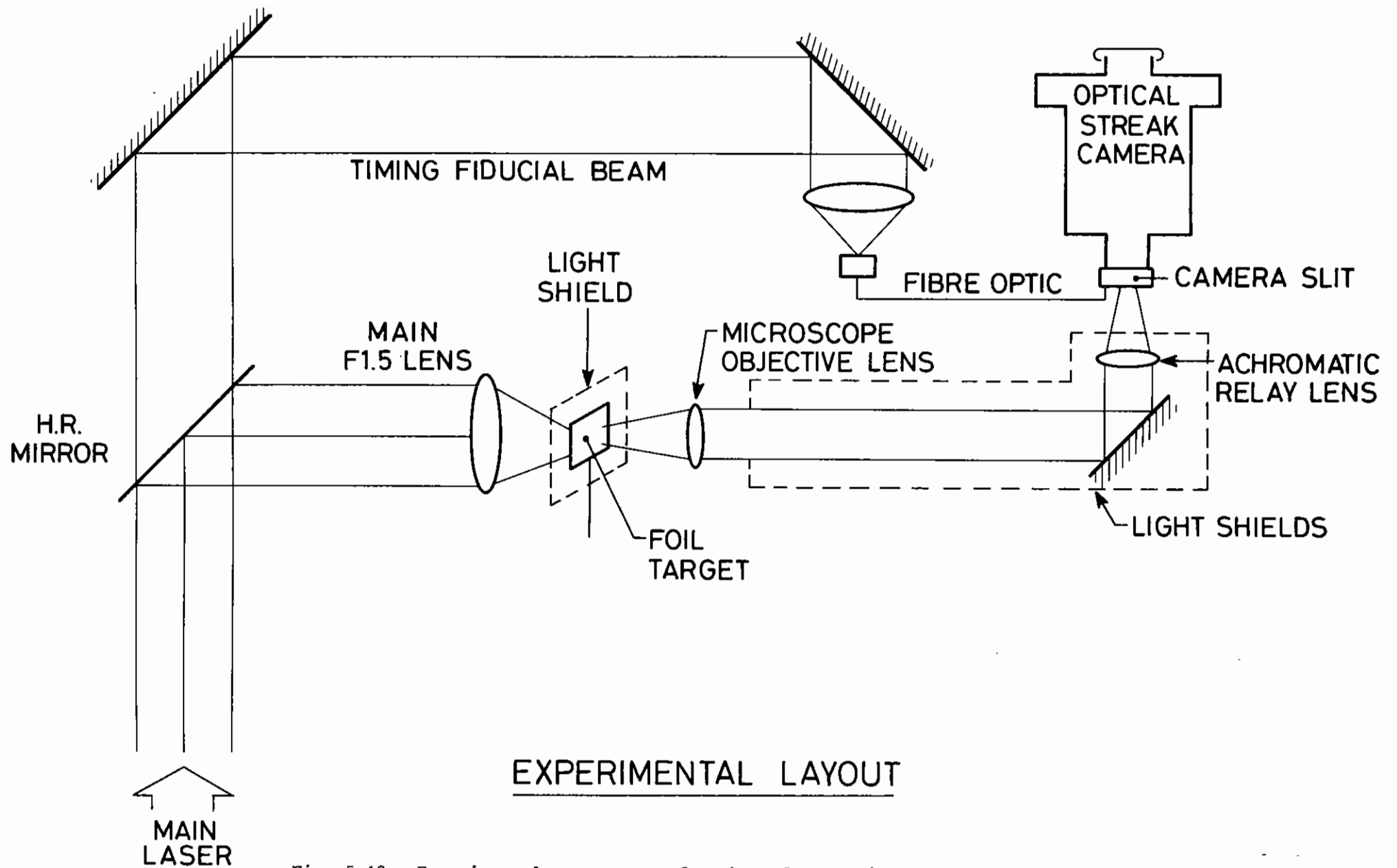


Fig. 5.19 Experimental arrangement for thermal smoothing and shock velocity measurements.

focused beyond the plane of the target to give a spot $\sim 150 \mu\text{m}$ diameter at the surface. A spatial period of $70 \mu\text{m}$ and intensity modulation ratio of 2:1 was measured from microdensitometer tracings of equivalent plane photographs recorded on type 55 polaroid film as shown in Figure 5.20. The targets were $25 \mu\text{m}$ thick Aluminium foils a few mm wide mounted onto a light shield to inhibit scattered laser light. The rear surface of the targets were viewed with a 25 mm 0.1 NA microscope objective lens and re-imaged with an achromatic relay lens onto a $100 \mu\text{m}$ slit on an optical streak camera fitted with a S20 photocathode.

The static spatial resolution was estimated to be $10 \mu\text{m}$ and the temporal resolution 10 ps .

The streak photographs for the unmodulated and modulated beam profiles are shown in Figure 5.21. The unmodulated image profile shows planarity corresponding to the unmodulated beam densitometer tracing (not shown here) limited by the spatial fill factor of 0.5 in the $f/1.5$ focusing lens.

It is interesting to note that there is no indication of any precursor emission occurring earlier in time before the emergence of the shockfront, which shows a risetime of less than 30 ps . The structure observable on the timing fiducial image is due to the position of the fibre optic well inside the focus of the fiducial beam lens. The lack of any observable structure in the unmodulated beam case negates the significance of fine scale modulations from inhomogeneities in the $f/1.5$ focusing lens, which are known to exist.

The degree of smoothing is expressed by the parameter δ where

$$\delta = \frac{T_{\text{MAX}}}{T_{\text{MIN}}} - 1$$

T_{MIN} is the time taken by the shockfront produced by a maximum in the drive pressure to propagate through the target and T_{MAX} is the time taken by the shockfront produced by a minimum in the drive pressure.

The variation of δ as a function of incident irradiance is plotted in Figure 5.22 and found to vary accordingly to the relation $\frac{1}{\delta} = 0.012 (I/10^{13} \text{ W cm}^{-2})^{1.3}$. There is little smoothing at low irradiance but smoothing increases at higher irradiance.

The small dataset over a limited region of irradiation makes the comparison of scaling with previous experimental results somewhat tentative. However it is of interest to note that the scaling from (5.29) is

$$\frac{1}{\delta} = 0.5 (I/10^{13} \text{ W cm}^{-2})(\lambda/1 \mu\text{m})^{1.5}$$

whereas from Gardner and Bodner (5.35)

$$\frac{1}{\delta} \propto I^{0.7} \lambda^{2.7}$$

If no smoothing was present one would expect the shock transit time to scale with the drive pressure by the simple inverse square law. Assuming the pressure varies as a function of irradiance according to $P \propto I^{0.7}$ (5.14) using the measured irradiance modulation gives a low irradiance limit of $\Sigma = 0.27$ to be compared with the observed value of 1.6.

This disagreement seems difficult to explain. Shock transit times are comparable to previous studies and although a wealth of processes exist to increase the observed smoothing parameter, anti-smoothing effects of sufficient significance are less abundant. Further analysis and experimentation is being performed to investigate this anomaly more fully. Although the zero smoothing value is somewhat questionable, a four-fold smoothing has been observed in the small irradiance region concerned. On this basis significant ten fold smoothing requires an irradiance on the limit of present laser performance. Further experiments will reveal whether this level of smoothing is achieved.

Despite the limited dataset obtained the experiment has demonstrated the

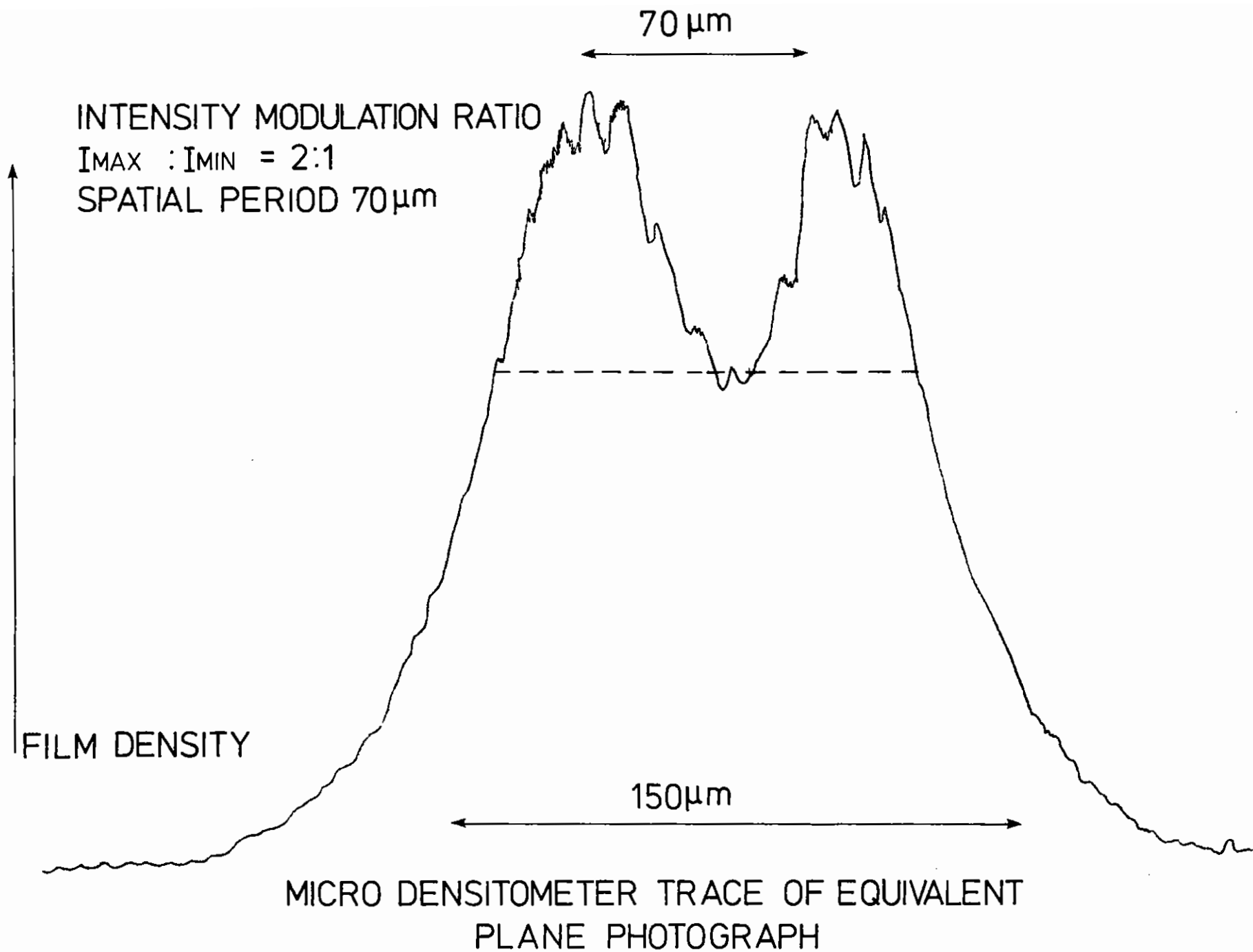
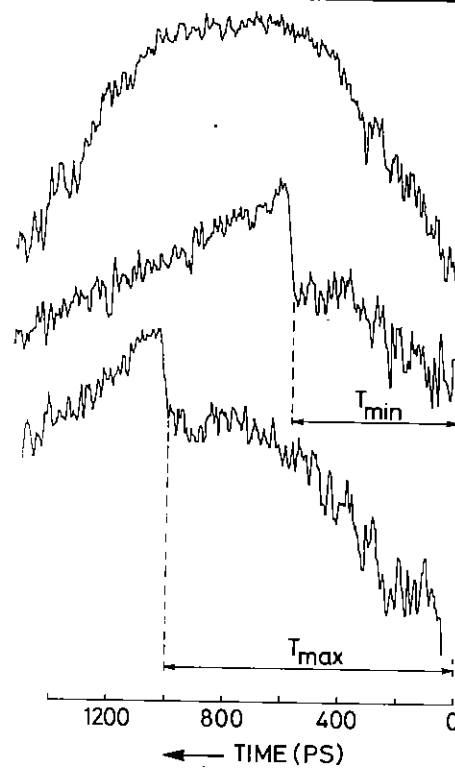
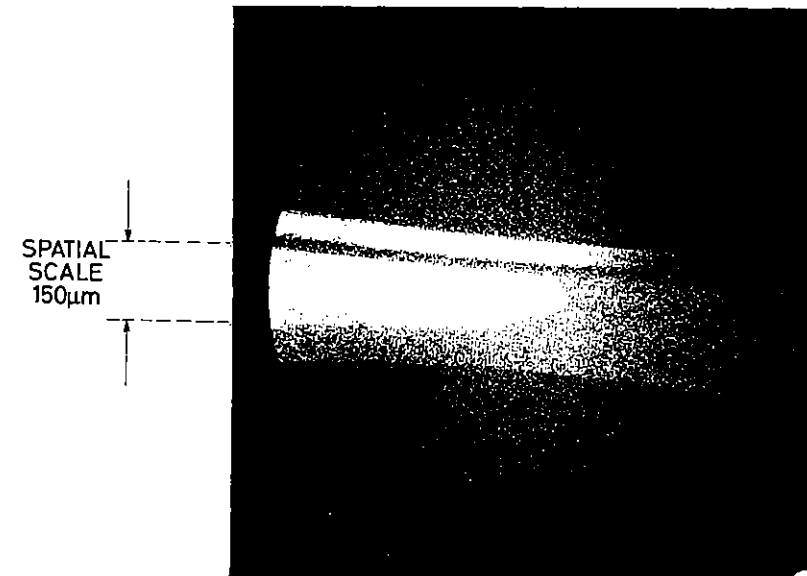
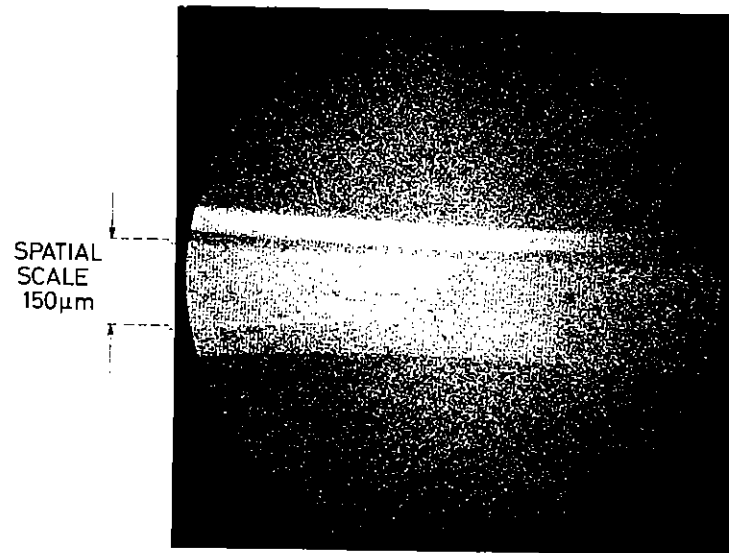


Fig. 5.20 Microdensitometer tracing of equivalent plane photograph showing imposed modulation of period $70\mu m$ and intensity ratio of 2:1. Dashed line represents half maximum.

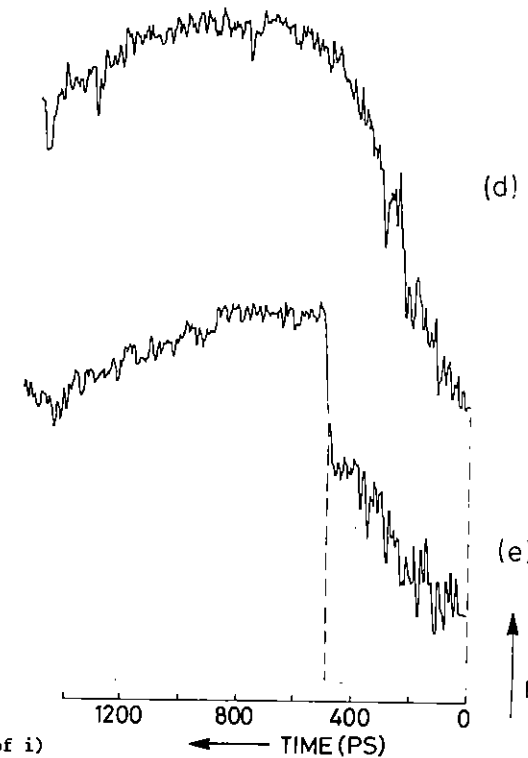


(a)

(b)

(c)

FILM
DENSITY

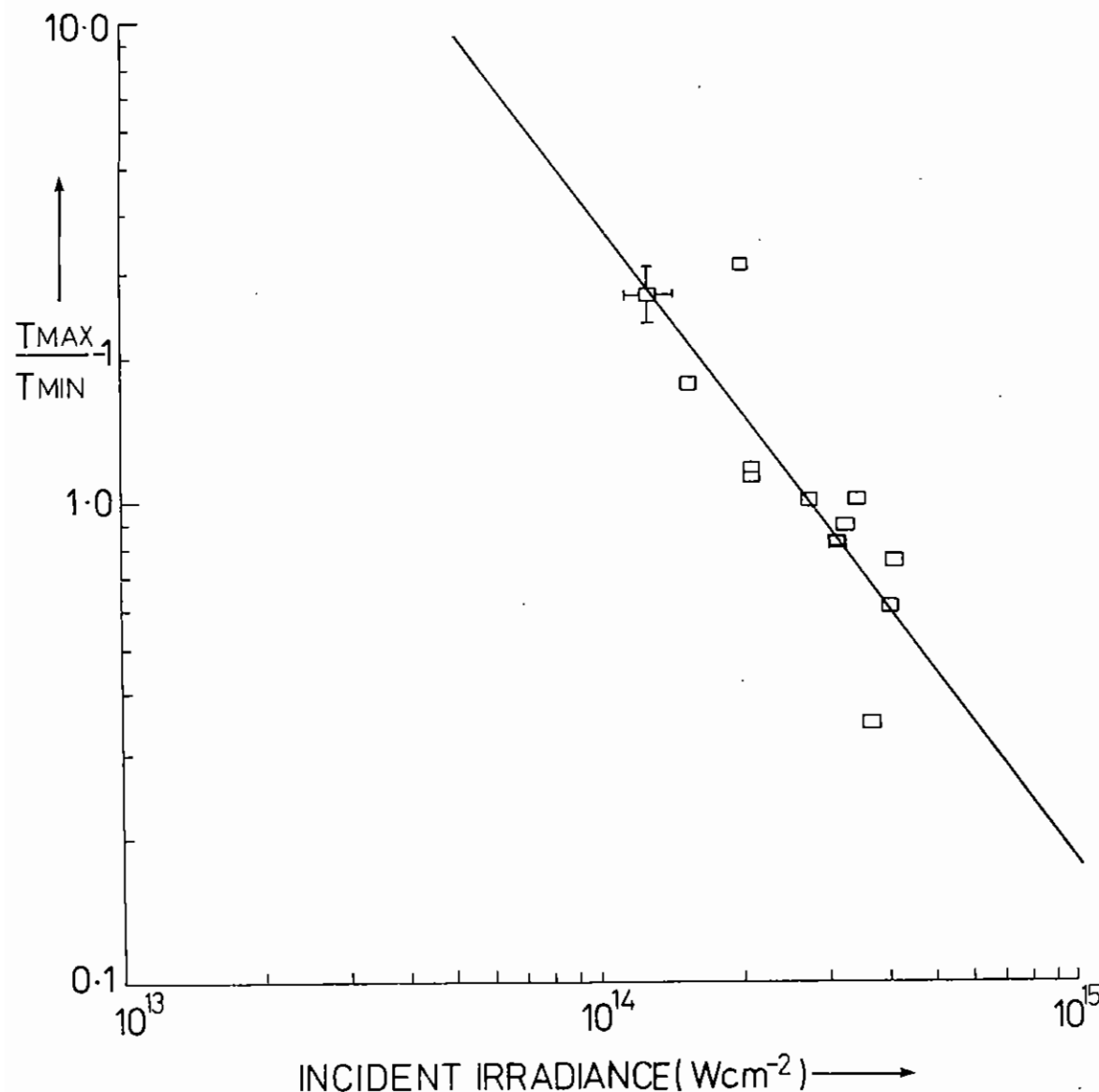


(d)

(e)

FILM
DENSITY

Fig. 5.21 Streak photographs of thermal emission from the rear surface of $25\mu\text{m}$ Al foils for the case of i) modulated laser beam, spatial period and intensity modulation as Fig. 2 and ii) unmodulated beam. Together with densitometer tracings of a) and d) laser pulse timing fiducial, b) and c) emergence of shockfront from rear of target at different times corresponding to pressure modulation at ablation front and e) unmodulated shockfront profile.



THERMAL SMOOTHING DATA AT $\lambda_L = 0.35 \mu m$.

Fig. 5.22 Variation of $\frac{T_{MAX}}{T_{MIN}} - 1$ as a function of incident irradiance for $\lambda_L = 0.35 \mu m$.

ability to measure smoothing parameters on a sub-nanosecond timescale to compare with previous longer time studies. Further experimentation and computer simulations are being performed which will provide a complementary dataset and an attempt will be made to directly quantify a relationship between the previous backlighting technique studies and the modulated shockfront technique. Of particular significance will be how the degree of smoothing varies as a function of target thickness at constant irradiance probing any time dependent smoothing processes.

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5.8 A measurement of the Stopping Power of Hot Dense Matter

The rate at which fast, heavy, charged particles loose energy in materials at ordinary temperatures has been extensively studied both theoretically and experimentally (5.39). Much less work has been done on energy loss in hot ionised media where the collisions are with free electrons so that the ionisation potential no longer sets a lower bound to the energy loss in a single collision and the total energy loss rate can be substantially enhanced (by as much as a factor of two (5.40)), provided the velocities of the free electrons are much less than the velocity of the heavy charged particle (the test particle). At still higher temperatures where the electron thermal velocity approaches or exceeds the velocity of the test particle the collisional loss rate decreases and can reach values much below the rate in the corresponding cold material. There are also many detailed considerations (5.40)(5.41) including the ionisation balance in the medium and the effective ionisation potential of partly stripped atoms, the effective charge of the test particle, the possibility of energy loss to collective modes in the plasma and so on. Interest in this topic has been raised by the possibility that beams of energetic ions may be used to heat and ablate the surface of a target containing D-T fuel, compressing the fuel and raising it to thermonuclear temperatures. An accurate knowledge of the stopping power of the hot ablating surface is essential for the design of such schemes and the theoretical models should be verified experimentally.

This is not easy. At a temperature of 100eV and a density of 1gm/cm^3 the energy density is $\sim 5 \times 10^6$ Joules/cm³ and any material heated to this temperature disassembles with the thermal velocity of the ions. There are few ways to create such energy densities in the laboratory and still fewer simultaneously to produce a sufficient flux of test particles to make a measurement of energy loss. Only one previous experiment has been reported (5.42). In it a 1 Mev deuteron beam with a current density at peak power of 50 KAc m^{-2} was used to heat a sandwich target whose rear surface was deuterated to a temperature ~ 40 eV. The residual energy of the deuterons at the depth of the deuterated layer was obtained from time of flight measurements on forward produced neutrons from the reaction $D + D \rightarrow \text{He}^3 + n$. Broad agreement with a model calculation was obtained. The temperatures and densities in such an experiment are limited by the peak currents available from pulsed diodes and only one type of projectile is suitable.

In this paper we give the first results of an experiment which is capable of obtaining data in a wider range of conditions and at higher energy densities. The arrangement is shown in Figure 5.23. The test particles were 3.5 Mev α 's from $D + T \rightarrow n + \alpha$ reactions in a laser-driven exploding pusher implosion. In the present experiment 1.5×10^6 α 's were produced having an instantaneous flux of $10^{22} (\text{cm}^{-2} \text{ sec}^{-1} \text{ steradian}^{-1})$ from a source of 20 μm diameter. Four of the six beams of the VULCAN laser (5.23) were used to produce this implosion. (The method is also applicable to 3 MeV protons from $D + D \rightarrow p + T$.) The other two beams were used to heat a small region of an energy loss foil placed at a distance of 3mm from the implosion to a temperature ~ 100 eV. The heating pulse was timed by introducing delays into the paths of the heating beams, so that its peak coincided with the arrival of the α particles at the foil. Since all six beams were derived from the same pulse from a master oscillator, there was no timing jitter other than might be due to the implosion dynamics. Those particles which had traversed the heated area and others which had passed through nearby cold regions of the same foil were registered in a CR-39 plastic detector 4cms from the source.

The excellent space and energy resolving properties of this detector(5.43)

permit measurements of the energy spectrum of each group of particles so that the change in dE/dX due to the heating can be directly observed. In the present experiment, multiple scattering in the energy loss foil led to some mixing at the position of the detector. However, a pure sample of "cold" tracks could be obtained and used to unfold the mixed distribution in the central region where "hot" tracks were concentrated.

It was found necessary to protect the surface of the CR-39 from bombardment by plasma particles with a 3.5 μm layer of mylar in order to obtain tracks of reasonable quality. (In a subsequent experiment the detector was placed further from the target, magnets were used to sweep plasma particles away from the detector and further improvements were made in the quality of the CR39 with the result that there was no observable surface damage.)

Figure 5.24 illustrates the method used to align the source, foil and detector. The D-T filled microballoon was first aligned at the focus of the four beams. A microballoon glued to the edge of the energy loss foil was then positioned via an auxiliary laser and microscope so that it was in the direct path from the centre of the detector to the α source. The heating beams were focused on the microballoon by observing the obscuration of each alignment beam on a telemicroscope sighting through the opposing lens. The foil was then raised into position by centering its shadow on the microscope image of the α source microballoon. The diameter of the foil microballoon determined the size of the heated area. The energy loss foil gave additional shielding to the CR-39 surface from plasma damage and the shadow cast by the foil can be seen on the etched material when it is suitably illuminated. A photograph is shown in Figure 5.25. The location of this shadow with respect to the fiducial marks on the detector gives direct confirmation that the alignment procedure was accurate.

The energy loss foil was 3.5 μm Mylar, inclined at 50° to present a 4.6 μm line of flight path to the α particles. This was judged to leave just sufficient residual range for effective measurement of the α particles in the CR-39. The laser beam axis was at 50° to the normal to the surface of

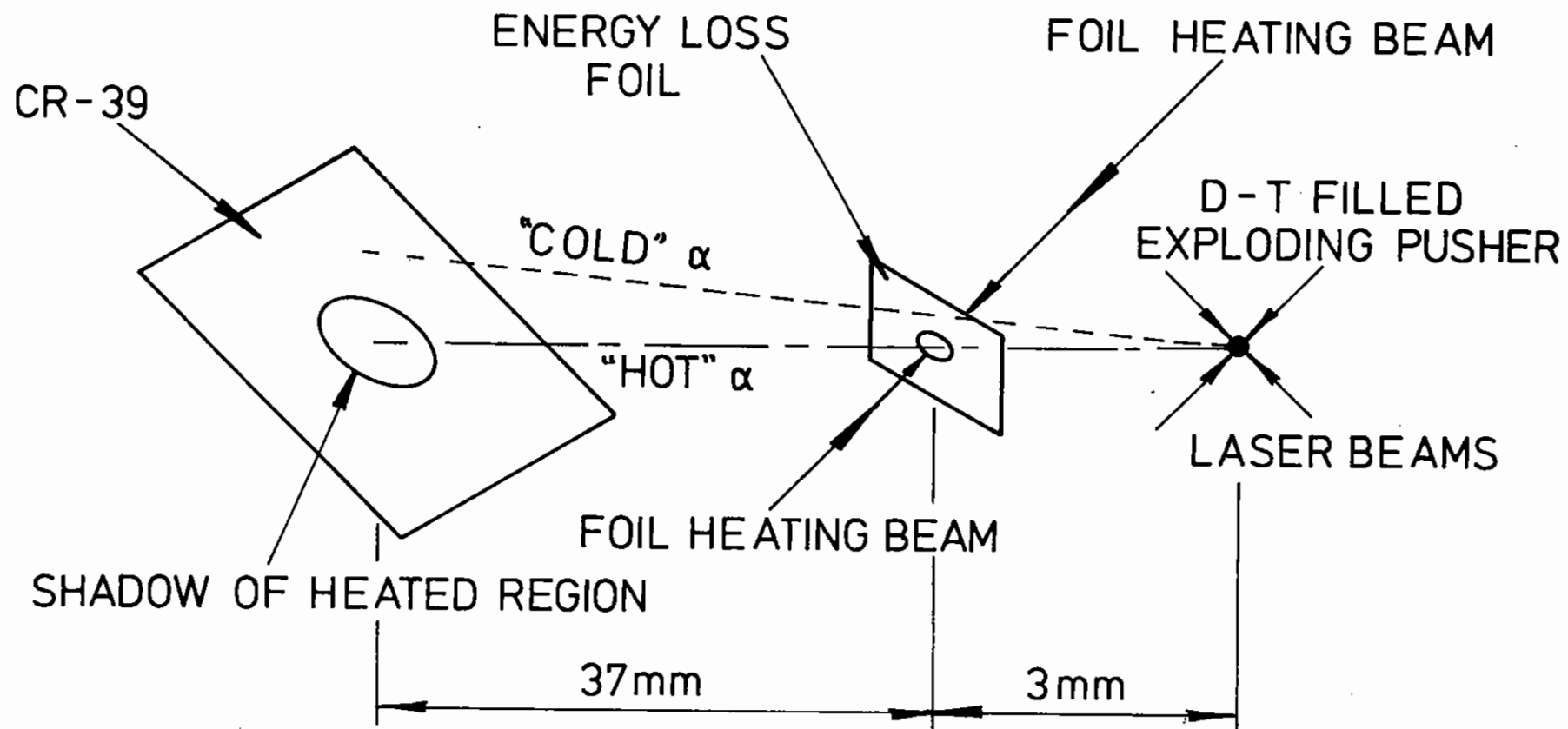


Fig. 5.23 Schematic design of experimental arrangement.

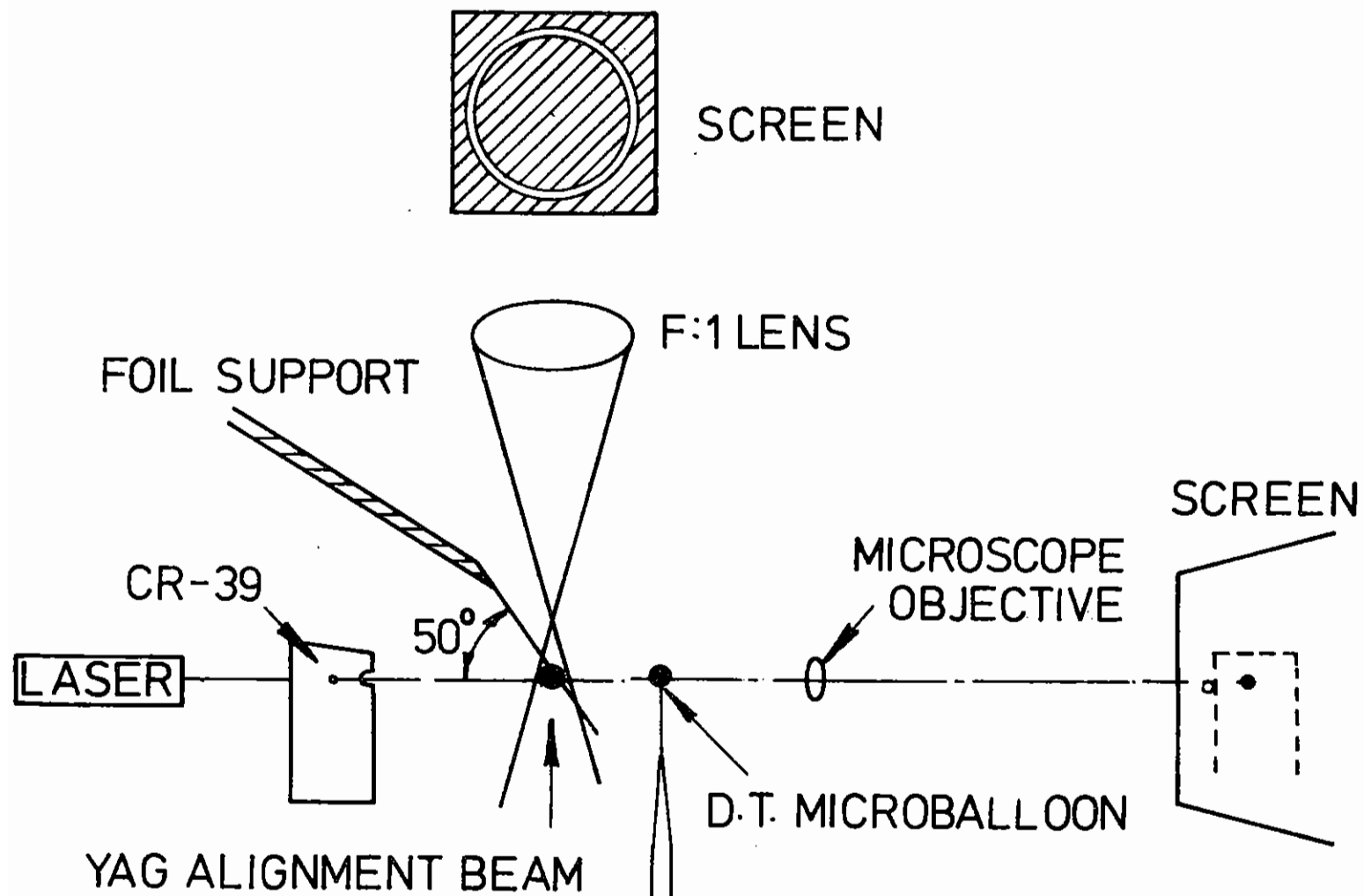


Fig. 5.24 Alignment method.

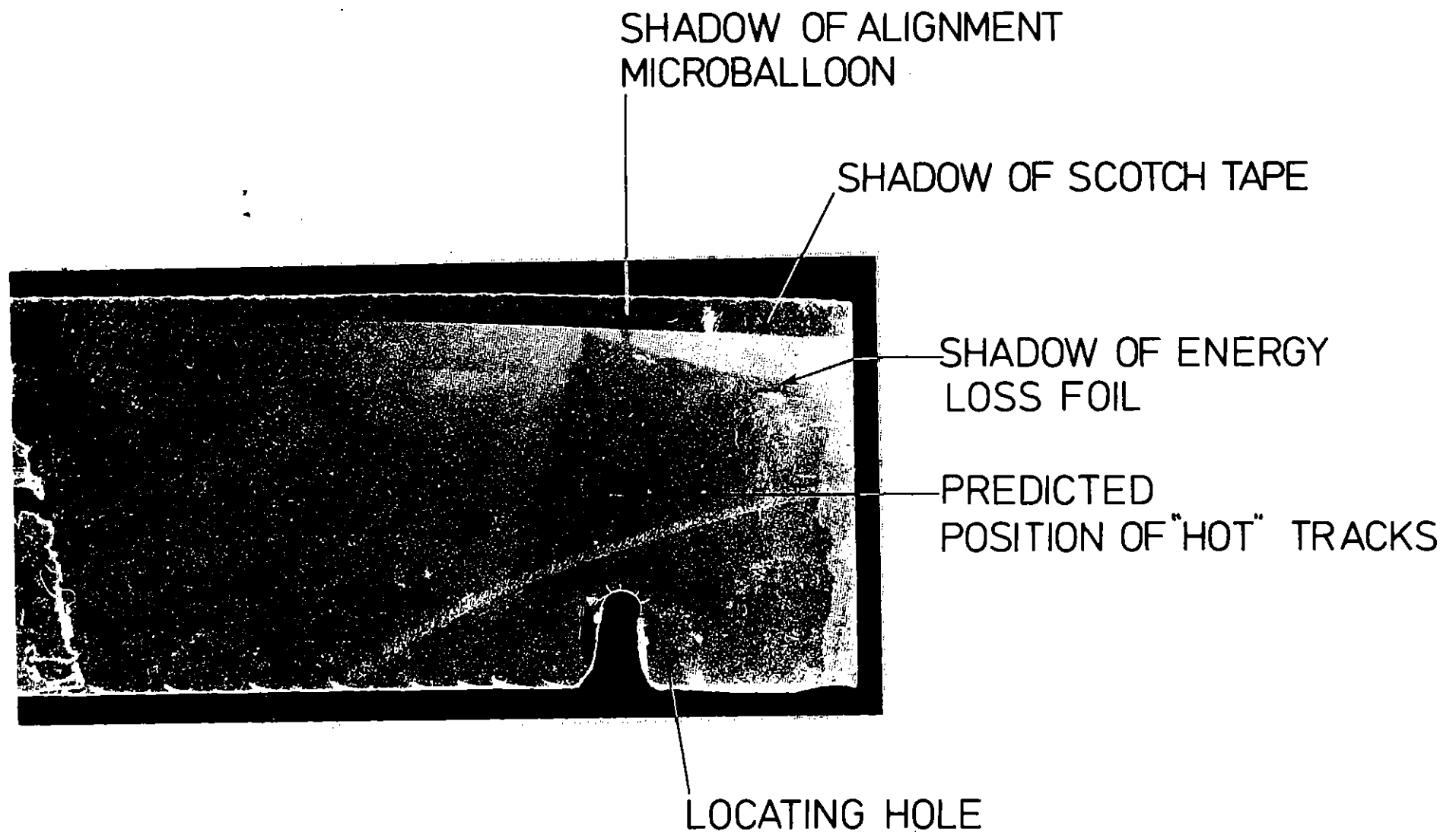


Fig. 5.25 Photograph of plasma damage on CR39 detector with illumination chosen to enhance visibility.

the energy loss foil. The two opposed converging beams do not completely overlap (as is shown in Figure 5.26(a)). The intensity distribution at the foil, projected on to the CR-39 is shown in Figure 5.26(b), and Figure 5.26(c) indicates the regions in which tracks were measured.

Optical streak cameras were used to set the time of arrival of all six beams. The implosion time of a similar microballoon was measured during the experimental run using an X-ray streak camera (5.44) to record both the prompt X-rays produced when the laser beams strike the target surface and the brief, delayed burst of X-rays at the peak of the implosion. The time differential was 150 ± 15 ps. In this type of implosion alpha particle emission occurs very close to the peak of the implosion. A 400 psec delay was introduced into both heating beams so that 3.5 Mev α particles would arrive at the foil in time with the peak of the heating pulse. Neutron yields were measured with a plastic scintillator detector placed 2 m from the target, outside the vacuum chamber. No attempt was made during this run to diagnose the behaviour of the foil. The greater part of the run was devoted to establishing the conditions under which good quality tracks could be obtained in the CR-39. One shot was successful in all respects. The shot parameters are given in Table 5.06.

The CR-39 polymer detector is manufactured under carefully controlled conditions at Bristol. After exposure it was etched in 6.25N NaOH at 60°C , $6.52\mu\text{m}$ of bulk surface being removed. These conditions were chosen to provide optimum energy resolution for tracks of range 5.5 to $7.0\mu\text{m}$ and dip angle 55° . The alpha particle tracks are clearly visible to the practised eye and track selection can be confirmed by measuring the angle that the major axis of the surface etch pit makes with the y-axis. The α -particle track density gave a yield of $(1.8 \pm .2) \times 10^6$ from the implosion which is in good agreement with the neutron yield of $(1.5 \pm .2) \times 10^6$ and confirms that the track observation efficiency is high even in the presence of plasma damage. Measureable parameters of an etched track are shown in Figure 5.27. In this experiment the etch pit minor diameter, M_1 , is the most sensitive measure of the residual range of the α particle and the presence of a change in residual range with position across the foil can be seen clearly in the raw data for the mean value of

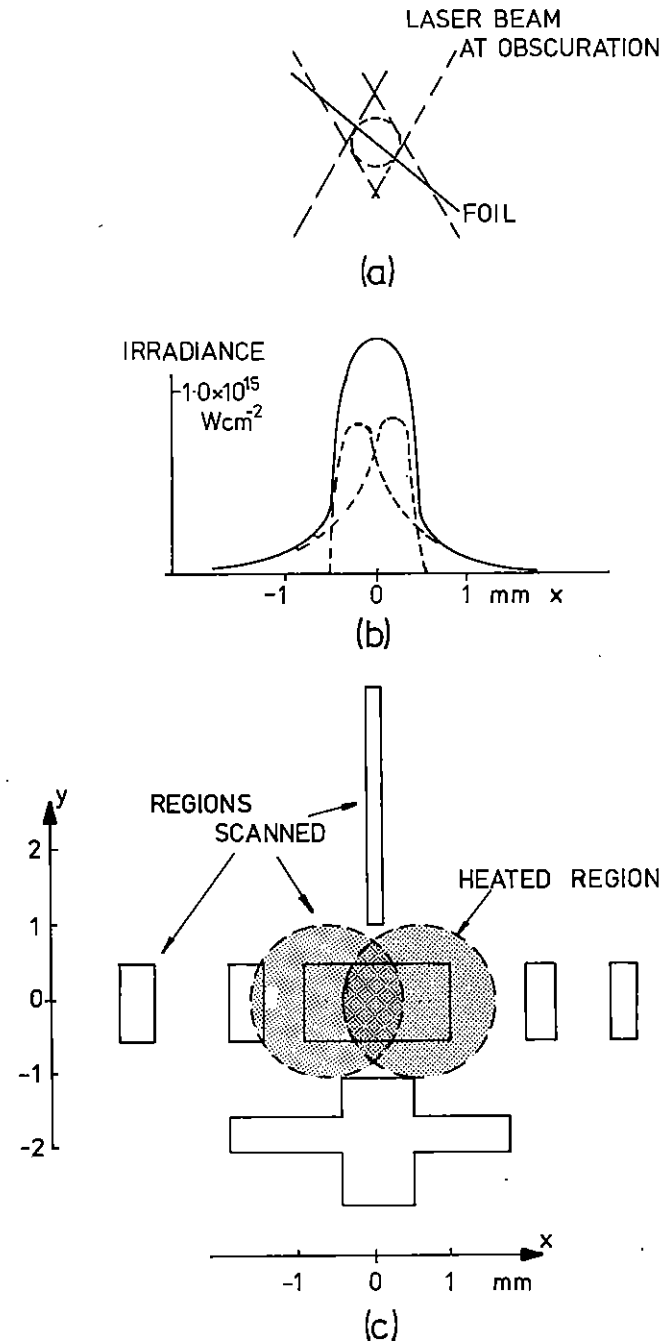


Fig. 5.26 a) Geometry of foil and illuminating beams. The dashed circle represents the microballoon used in the alignment procedure.
b) Horizontal intensity distribution at foil, projected onto CR39 detector. The dashed curves are for the individual beams. The solid curve is the sum.
c) Map of foil showing regions scanned and expected shadow of heated region.

Table 5.06 Shot Parameters

Total Energy of 4 beams:	40 J
Total Energy of heating beams:	20 J
Pulse Length:	100 ps
DT Target Diameter:	77 μm
Wall Thickness:	1 μm
Fill:	20 Atmos equimolar DT
Neutron Yield:	$(1.5 \pm 0.2) \times 10^6$
α -particle Yield:	$(1.8 \pm 0.2) \times 10^6$
Heating Beam Diameter:	100 μm nominal*
Energy Loss Foil:	3.5 μm Mylar
Orientation of normal to surface relative to heating beams:	50°
Orientation of normal to surface relative to α flight path:	50°

* See text.

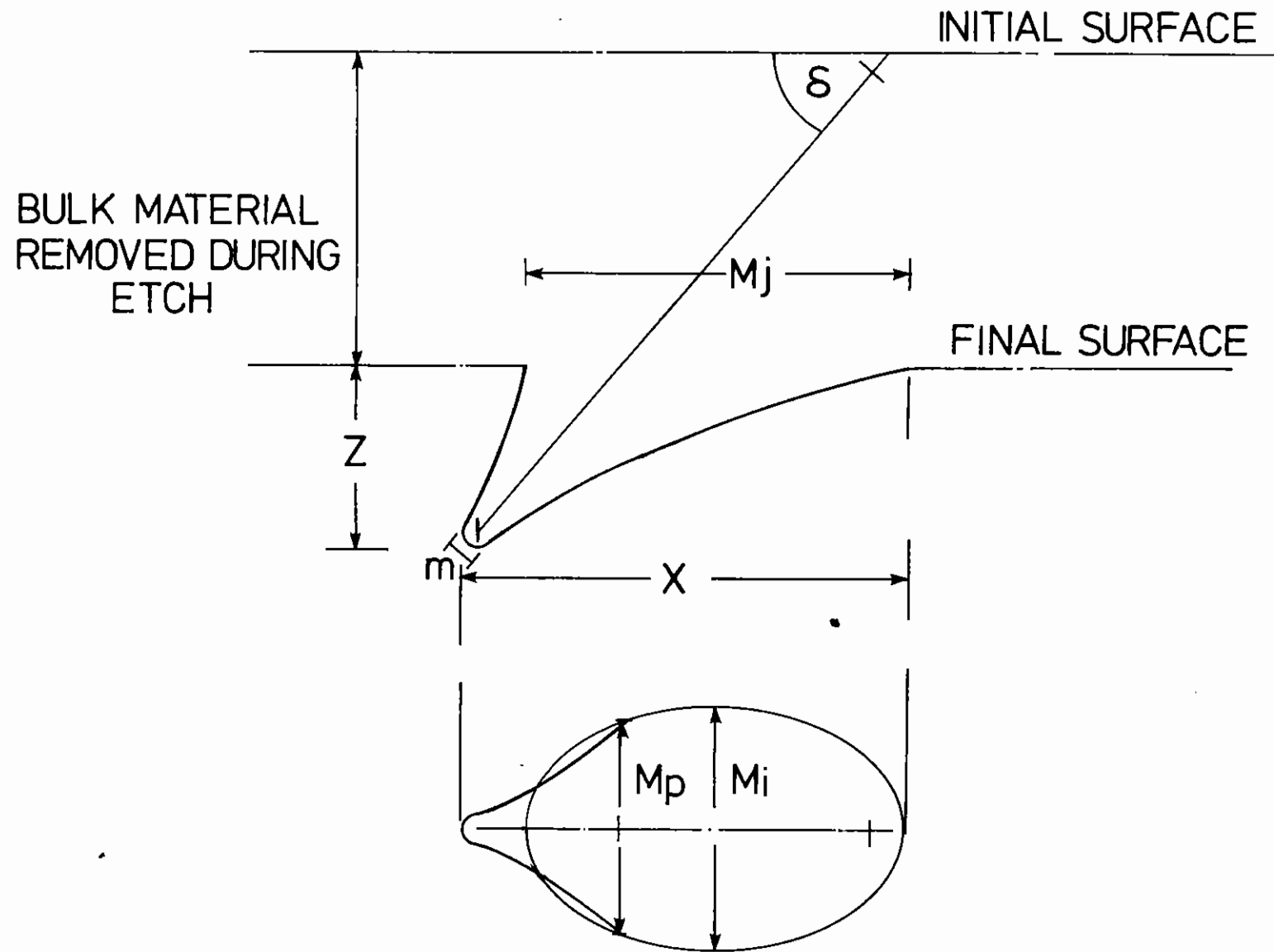


Fig. 5.27 A track simulation showing the five measurable etch-track parameters Z , X , M_j , M_i and m . Note that M_p , the diameter of the mouth where the track trajectory intersects the post-etched surface, is not a measurable parameter.

M_1 as a function of position shown plotted in Figure 5.28. A gradual linear variation due to the variation of dip angle with position has been subtracted from the data in Fig 5.28(a). The peak in M_1 indicates a reduction of residual range and it is centred within 0.1 mm of the expected position of the shadow of the hot region in the y-direction and within 1mm in the x direction. The y -distribution has a width (FWHM) of 1.1mm and the x-distribution has a width of 3.3mm. Coulomb scattering (computed from cold material radiation lengths) and the geometry of the irradiation with tilted foils suggest a width of 1.65mm in y and 3mm in x, although the temperature is expected to have some x dependence. The spatial distribution of the signal is thus in reasonable agreement with expectations. The scan could not usefully be extended further in the x direction. Unfortunately, the plasma damage was too severe to measure tracks in the region not shadowed by the energy loss foil.

The structures of the etched particle tracks in CR-39 are uniquely determined by the range and dip angle of the incident particles together with the variation of the track etch rate (V_T) with the residual range. Measurement of the parameters of the etch pit (Figure 5.27) together with suitable analysis can thus give a precise estimate of the range of the incident particle (5.43). It is first necessary to determine the variation of V_T with residual range. In this experiment it was convenient to produce a graph of the variation of the track parameter X as a function of M_1 and adjust the normalisation of the V_T -residual range relation until the computed trajectory fitted the measured points. The experimental X, M_1 points are shown in Figure 5.29, together with some theoretical trajectories. It is not necessary to know the residual ranges of the tracks used.

The best fit theoretical X - M_1 trajectory shown in Figure 5.29 has a V_T at 10 μ m residual range of 0.625 μ m hr⁻¹. The other curve shows the predicted trajectory where the whole V_T - residual range curve has been scaled to make V_T at 10 μ m equal to 0.750 μ m hr⁻¹ to indicate the sensitivity of the procedure. This V_T calibration procedure has been shown to be valid by using α -particles of known range from a Cf²⁵² source. From the V_T residual range curve thus obtained, the variation of M_1

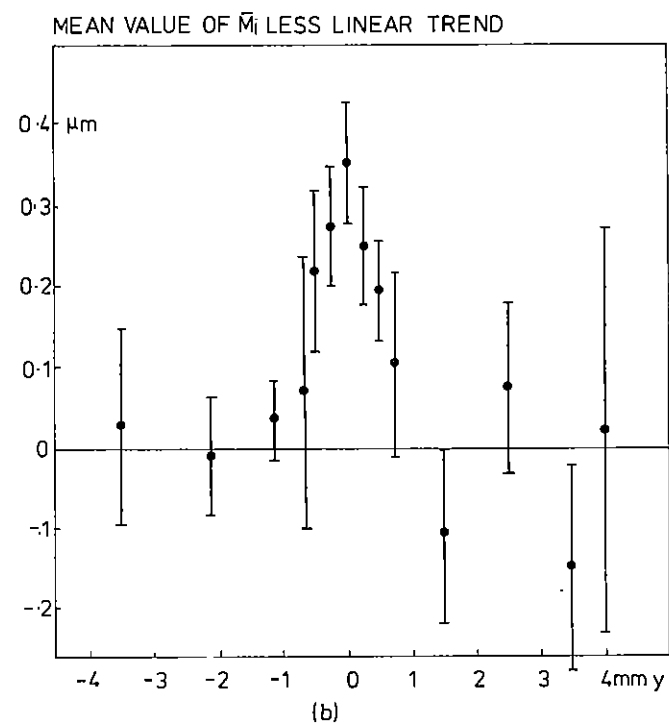
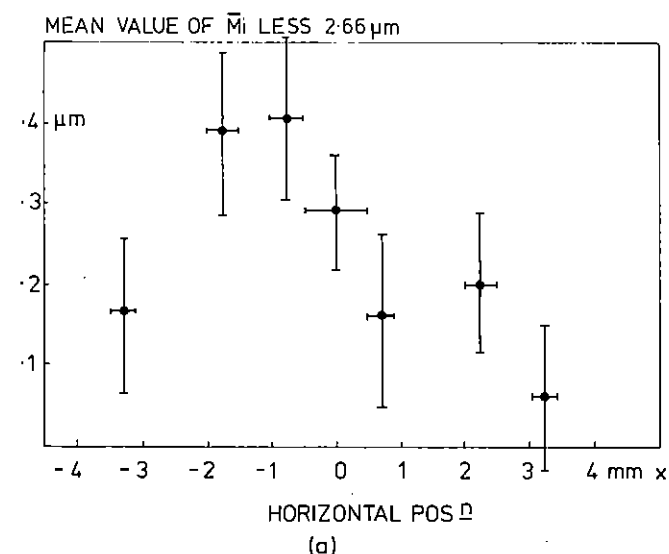


Fig. 5.28 a) Departure of mean value of M_1 from $M_1 = 2.66 \pm .14\mu$ as function of y coordinate for tracks with $|x| < 0.5$ mm. b) Departure of mean value of M_1 from 2.66 μ m as function of x coordinate for tracks with $|y| < 0.5$ mm.

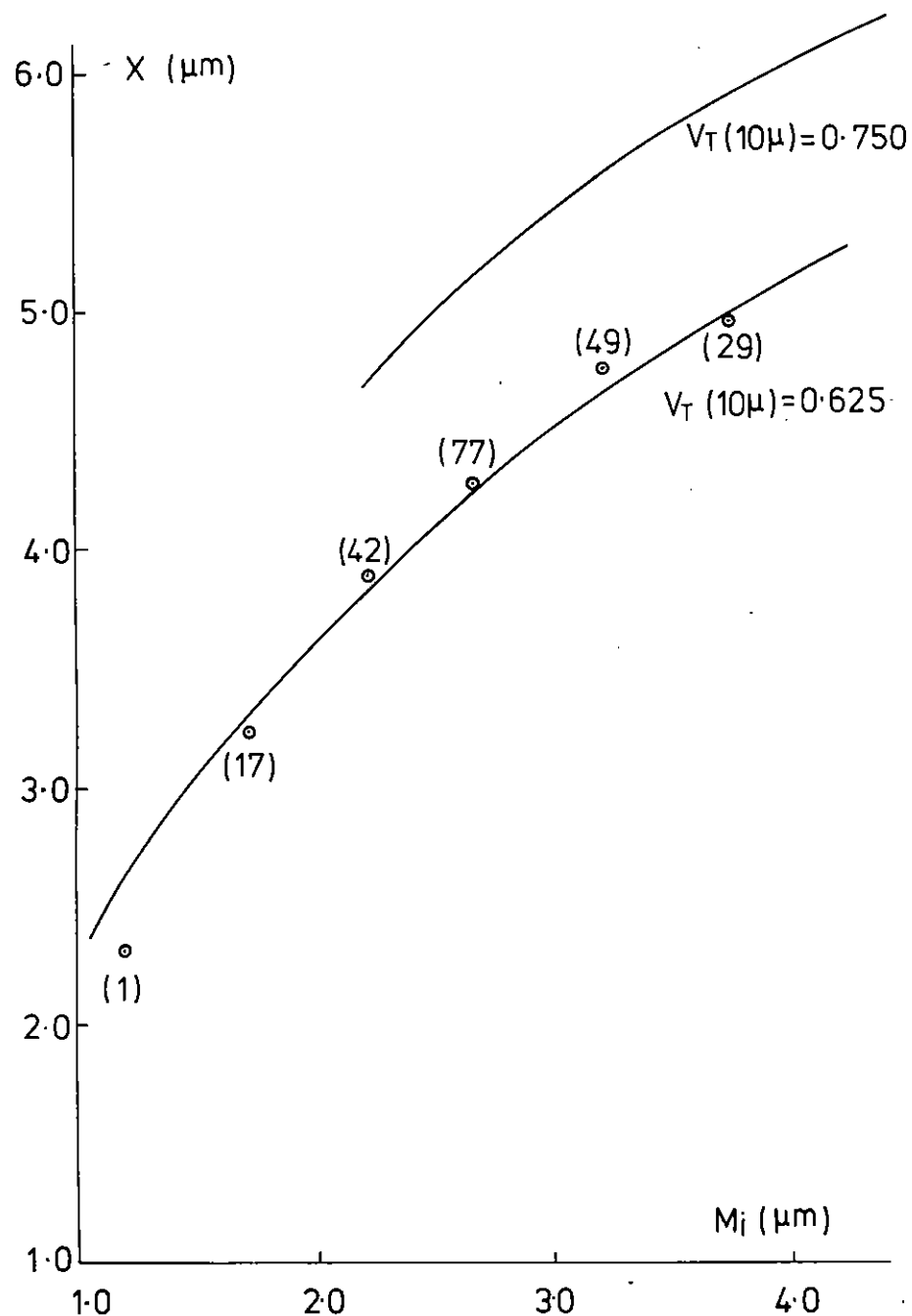


Fig. 5.29 Track parameter correlation x vs M_i for various etch rates.

with α -particle range was computed, enabling the energy of each incident α -particle to be determined. Figure 5.30 shows the variation of M_i with the incident particle range for a variety of dip angles used in the experiment.

We can make an immediate rough estimate of the peak value of the change in residual range due to the hot region by using an approximate estimate of Coulomb scattering dilution. At the position of the peak we expect a priori $\sim 1/3$ of all tracks to be "hot". Thus the difference in M_i between hot and cold is three times the height of the peak in Figure 5.28, or $1.3 \mu\text{m}$. This translates into a $3 \pm 1 \mu\text{m}$ difference in residual range, i.e. a 60 ± 25 percent change in foil stopping power. Figure 5.31(a) is a histogram of the residual range of tracks displaced in y by more than 1mm from the centre of the shadow of the heated region which the distribution of Figure 5.28(a) suggests should have no contamination of "hot" tracks. The mean residual range is $7.01 \pm .05 \mu\text{m}$ corresponding to an energy of 1.7 Mev .

When corrected for the intervening material, using cold material stopping powers, the α -particle range at source is $16.3 \mu\text{m}$ (CR-39 equivalent). The range of a 3.5 Mev α particle in CR-39 is $17.5 \mu\text{m}$. We attribute the difference to the stagnated pusher material. It is equivalent to $2 \times 10^{-4} \text{ gm cm}^{-2}$ of cold SiO_2 which is entirely consistent with exploding pusher behaviour. The width of the distribution is much greater than the resolution of the detector (which approaches the straggling limit (5.43)) and therefore represents the spread in α energy at source. This can be due to the source temperature (which gives a spread in CMS velocity in the lab), to any possible space or time dependence of pusher ρAR or to a time-dependent source electrostatic potential (5.46). Because the laser pulse length was less than the implosion time, we do not expect this last to be significant. If we assume the width to be wholly due to the source temperature, we obtain a value of $1.2 \pm .4 \text{ keV}$. This can be compared with a source temperature deduced via a simple uniform temperature and density model from the magnitude of the neutron yield of $1.2 \pm .3 \text{ keV}$.

In Figure 5.31(b) we show the corresponding histogram of tracks taken from

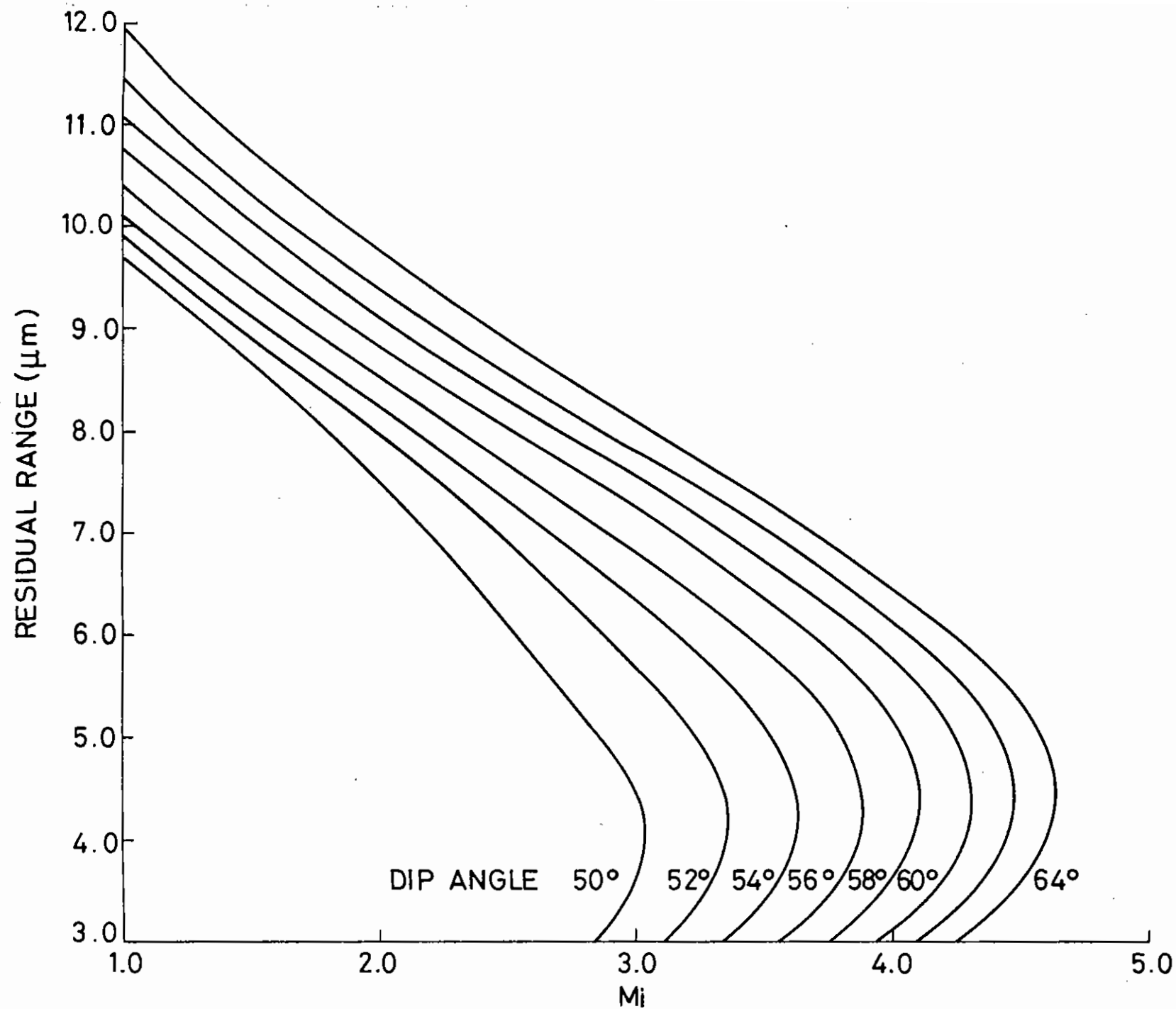


Fig. 5.30 Calculated residual range as function of track width for different dip angles.

the region $0 > x > 2\text{mm}$, $y < 0.5\text{mm}$ (see Figure 5.26(b)) and in Figure 5.31(c) a subtraction of Figure 5.31(a) from Figure 5.31(b) in which the number of tracks in Figure 5.31(a) has been normalised so that the region with residual range $> 7 \mu\text{m}$ contains the same number of tracks as the corresponding region of Figure 5.31(b). The subtracted distribution has a mean residual range of $5.45 \mu\text{m}$ CR-39. The change in stopping power due to heating is thus equivalent to $1.56^{+0.8}_{-0.4}$ of CR-39 or $1.49 \mu\text{m}$ mylar, i.e. there is $33^{+16}_{-8}\%$ enhancement. This is in agreement with the approximate estimate based on predicted coulomb scattering mixing. The spread in hot track residual energy is consistent with the value for cold tracks. This confirms the validity of the (nonlinear) conversion of track width to residual range and also gives us confidence that the change in stopping power is sufficiently uniform for 1-dimensional modelling to be meaningful.

Heating of the energy loss foil was modelled in 1-D using Medusa (5.47) assuming an absorption of 25% of which 10% is in the form of hot electrons with $T_{\text{hot}} = 10 \text{ keV}$. A half thickness foil with beam incident from one side was considered. The deposition of the hot electron energy was computed by the method of Harach and Kidder (5.48). The temporal history of the median plane is shown in Figure 5.32(a). Very early in the pulse there is significant shock compression followed by a rapid decrease in density due to the fast electron driven expansion. The temperature rises to a peak towards the end of the pulse. Figure 5.32(b) shows the spatial variation of the foil parameters at $t = 125 \text{ ps}$, corresponding to the peak of the heating pulse. The alpha particles traverse the foil at $t = 112 \pm 20 \text{ ps}$.

The free electron stopping power was computed using binary collision theory as quoted in (5.40) and is shown in Figure 5.32(c) for three time steps. The additional stopping power due to the remaining bound electrons is small. We estimate a 5% increase in the mean value at $t = 100 \text{ ps}$ and a less than 2% increase at $t = 125 \text{ ps}$. The mean stopping power passes through a maximum of $1560 \text{ MeV cm}^2 \text{ gm}^{-1}$ for 3.0 MeV alpha particles at $t = 112 \text{ ps}$ so that our result is insensitive to any small timing error.

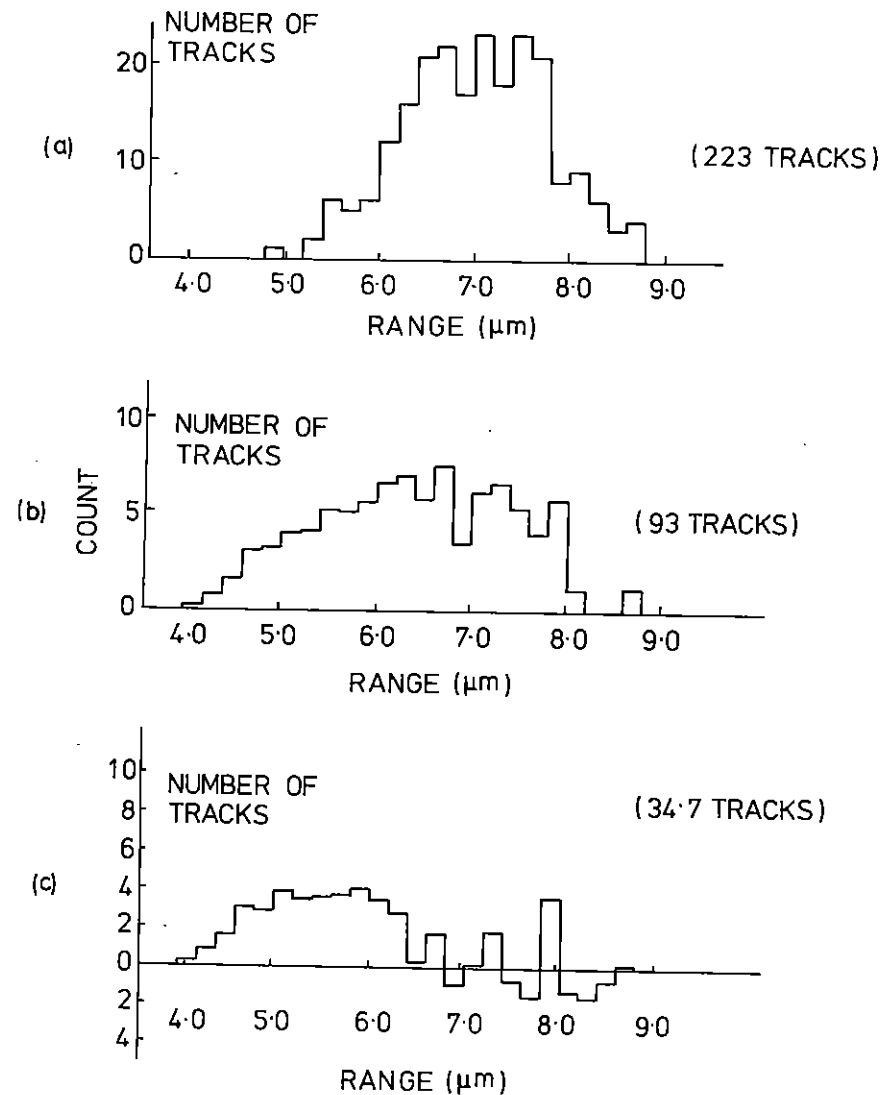


Fig. 5.31 Residual range distribution

- a) $|y| > 1\text{mm}$ "cold" tracks. Mean $7.01 \pm .05 \mu\text{m}$.
- b) $|y| < 0.5\text{mm}$, $|x + 1| < 1\text{mm}$: "mixed" tracks. Mean $6.41 \pm .10 \mu\text{m}$.
- c) Subtraction b) - a): "hot" tracks. Mean $5.45^{+0.4}_{-0.8} \mu\text{m}$.

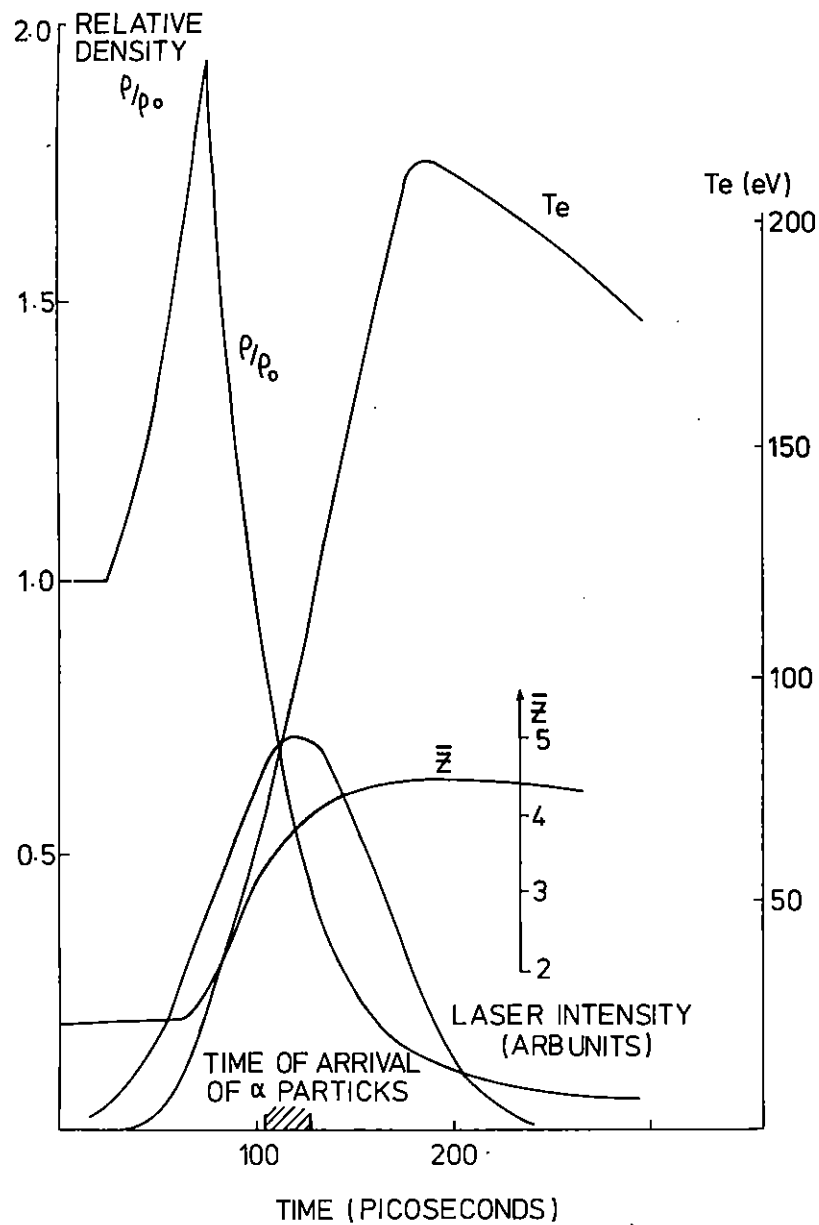


Fig. 5.32 a) Time history of foil parameters at centre of foil.

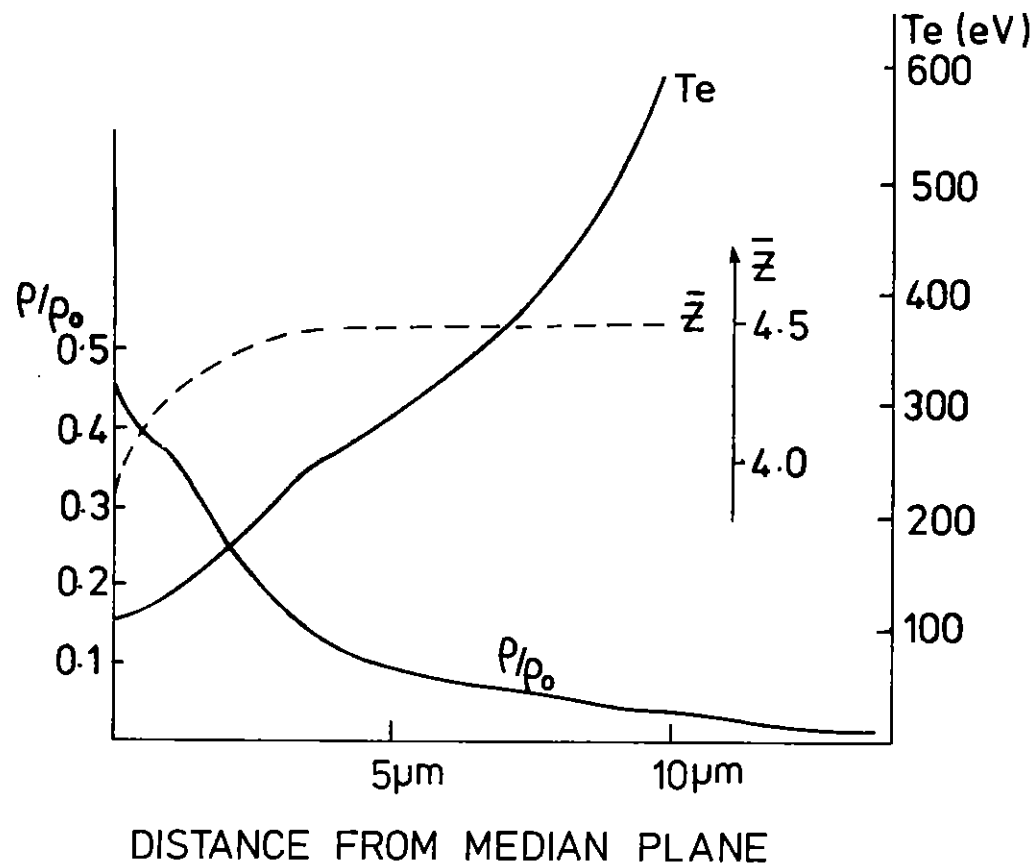


Fig. 5.32 b) Spatial variation of foil parameters at time of passage of α -particles.

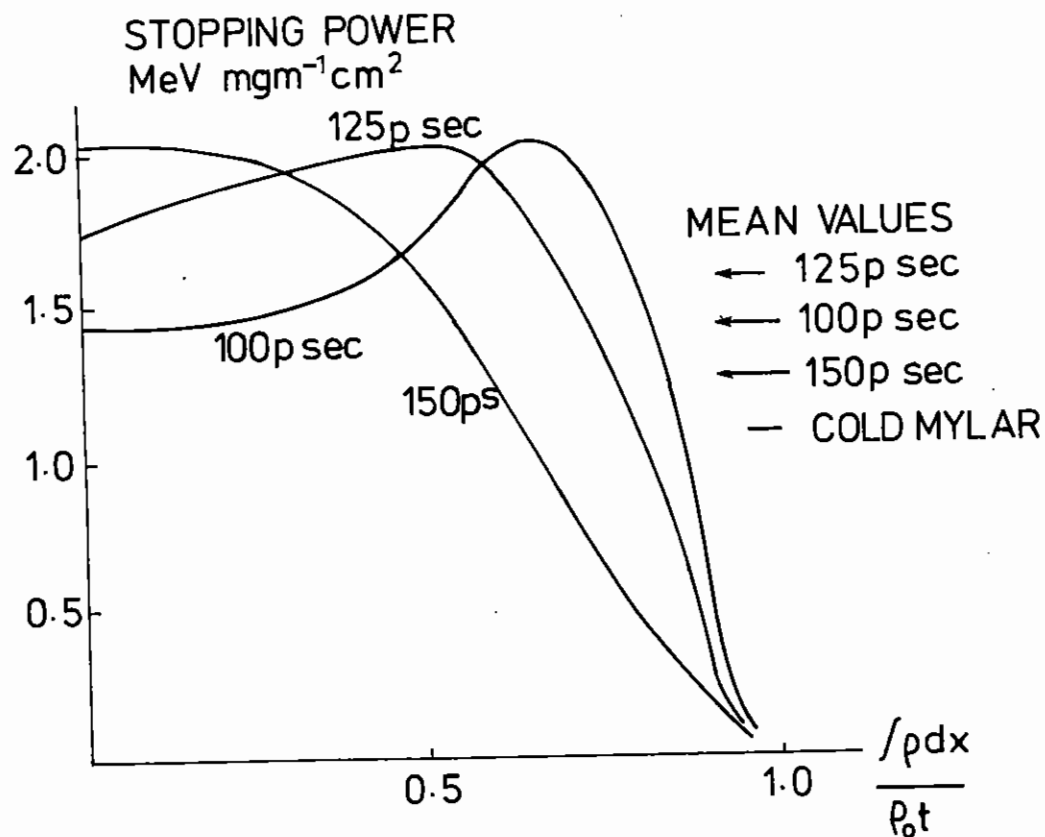


Fig. 5.32 c) Free electron stopping power of foil at three time intervals.

Because the low density plasma at large distances from the median plane of the foil contributes very little to the total stopping power we are insensitive to 2-dimensional effects in the plasma expansion. We have yet to study in detail the effect of lateral variations in the irradiation intensity. The results of Endo and Doke (5.49) and Ziegler (5.50) were used to compute the stopping power of the cold foil. Taking account of the variation of stopping power with energy we find a predicted ratio of "hot" to "cold" stopping power of 1.40 which is in excellent agreement with our experimental value of $1.33^{+0.16}_{-0.08}$, where the error quoted is purely statistical.

Further experiments are planned covering a range of materials and conditions.

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5.9 Laser Scattering Measurements of Heat Flow at High Temperature Gradients

In this experiment the thermal conductivity of a uniform field free plasma is measured under conditions where the electron mean free path λ_e is equal to the temperature gradient scale length, $L \equiv T / \nabla T$. A z-pinch is used to create a column of plasma at 4.5 eV some 3 cms in diameter and 25 cm long. The number density may be chosen to lie between 5 and $15 \times 10^{16} \text{ cm}^{-3}$, by varying the filling density of Hydrogen and the capacitor bank voltage. The number density and temperature remain constant for approximately 1 μs .

1 μm radiation in a 4 ns flat topped pulse is focused radially at $f/10$ onto the axis of the pinch to form a spot estimated to be 120 μm diameter. The evolving thermal wave-front is then followed using laser (Thomson) scattering at 90° of frequency doubled green light at 0.53 μm . Between 1. and 12 J of green light was focused at $f/10$ to a spot size estimated to be of similar dimensions to the infra-red beam. The scattered light was collected at $f/12$ and directed into a Spex spectrometer and then fed into

a Hadland type 675 streak camera. This allowed temporal resolution of the entire scattered spectrum.

In order to collect enough green scattered light to obtain sufficiently good statistics to follow the changing plasma parameters with a time resolution of about 500 ps, a slit of 200 μm was used at the Spectrometer input. It follows that the plasma parameters are measured along a cylindrical volume 200 μm long and 120 μm diameter at 45° to the heated volume. This is illustrated in Figure 5.33 which shows the interaction region at the axis of the pinch both from above and from the side.

The measured parameters are therefore averaged over this volume. To obtain information about the spatial behaviour of the thermal wavefront the infrared beam was moved horizontally and vertically. It was not possible to measure accurately the position of the infrared beam with respect of the scattering volume, although the amount by which the infrared beam was moved is known with more certainty. Thus it was necessary to scan the infrared beam until a temperature maximum was found and this was assumed to be the position of coincident foci of the two beams. The positions of the other measurements were then known.

The differential scattering vector, K_s , could be chosen to be either parallel or perpendicular to ∇T by having the green beam enter in either of two opposing directions. Both were tried successfully. Because the probe beam was derived from the infrared beam, it only lasted 4 ns. With a longer probe beam pulse it would be possible to follow the whole time-history of a particular event. For many shots the relative timing of the beams was such that the probe beam illuminated the plasma for 1 ns before the arrival of the heating beam, allowing the initial z-pinch parameters to be found.

It was not in fact possible to measure the unperturbed initial temperature because the probe beam was sufficiently intense to heat the plasma up to 8-9 eV in 4 ns (for a 10 J pulse, assuming Spitzer conductivity). The initial temperature was assumed to be 4.5 eV from previous use of the pinch. This means it is not possible to measure the temperature change

due to the heating beam at distances greater than about 1 mm from the beam axis as the temperature does not exceed 8 - 10 eV at any time in this region, but this does not pose a serious limitation in interpreting the results.

Results

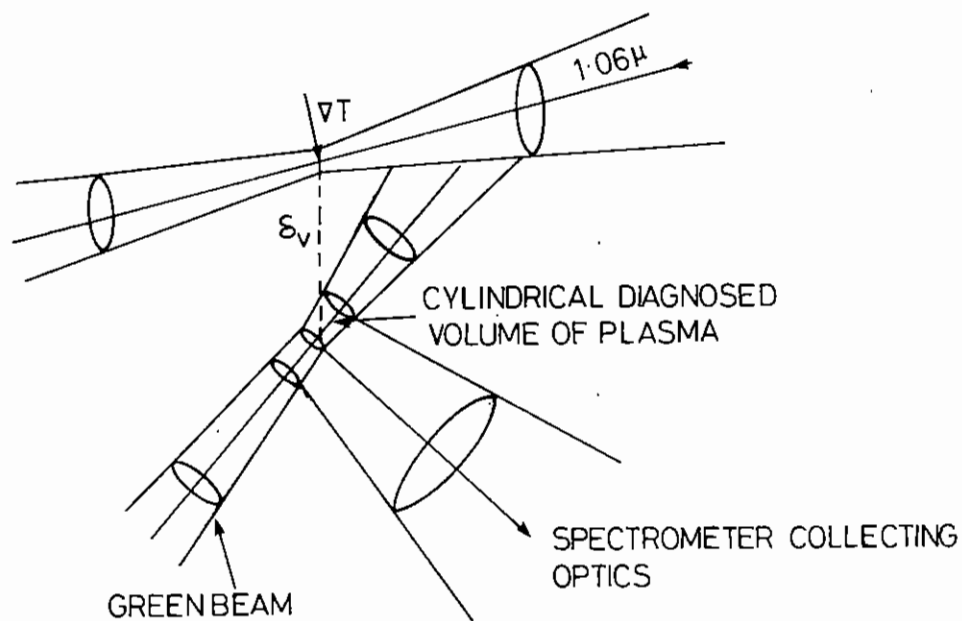
An example of a typical streak with and without the heating beam is shown in Figure 5.34, together with the corresponding microdensitometer traces at one 500 ps sampling interval in each trace.

The traces show the broad electron feature which is spectrally resolved and the central ion feature which is integrated over. The resolution of the diagnostic was approximately 7 $\text{\AA}$. In most shots where there is enough light in the electron feature to gain a good value for N_e and T_e , however, the ion feature saturates the film. The ion feature is enhanced in nearly all shots where the plasma is heated by the 1 μm beam, but because of the saturation of the film it is only possible to put a lower bound on the degree of enhancement. This lower limit ranges between 2 and 7, and applies to the situation of K_s perpendicular and K_s parallel to ∇T . This is very similar to the results of Grey and Kilkenny (5.53) and is taken as direct evidence of ion-acoustic turbulence generated by the return current associated with the heat flux.

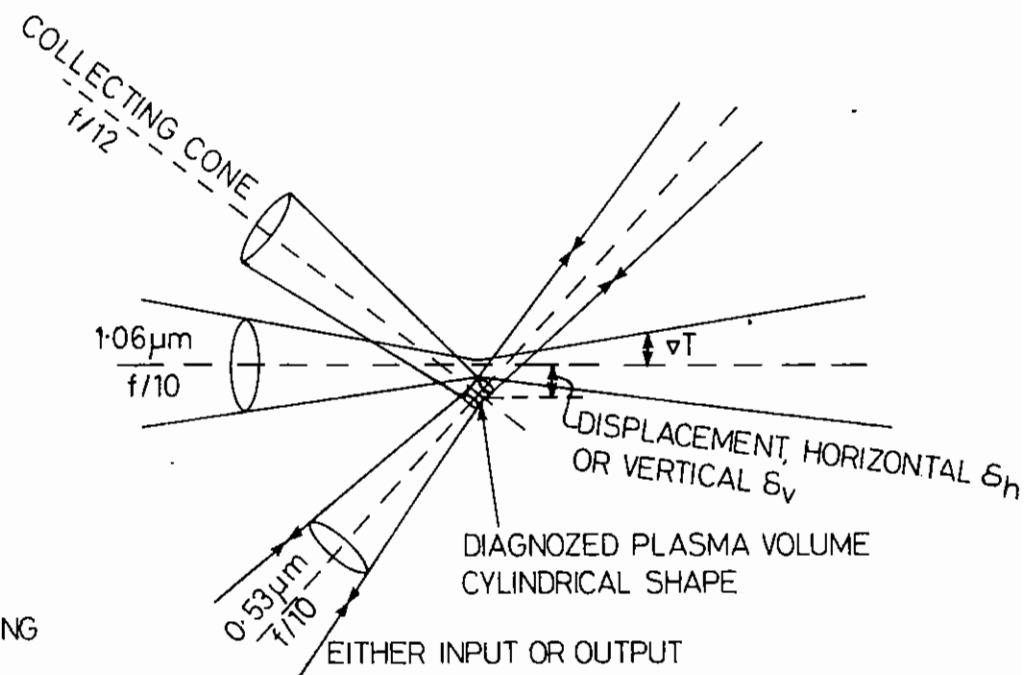
The analysis of the electron feature yields a value for the density and electron temperature. This analysis follows the method devised by Kunze (5.52). In Figure 5.34 we see the value of α fall, corresponding to a temperature increase as the plasma is heated. The two values for each spectrum correspond to approximately 14 eV and 8 eV.

In a significant number of shots an asymmetry is observed in the shape of the spectrum of the electron feature, the electron feature yielding a different value of α (and hence N_e and T_e) for each side of the spectrum.

This asymmetry occurs particularly at the times of greatest ∇T . The



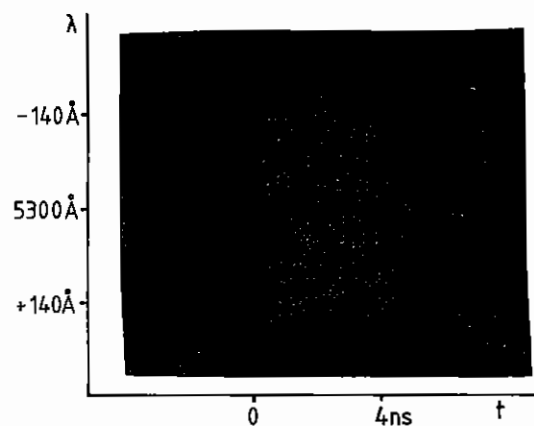
A RADIAL VIEW OF THE INTERACTION REGION.



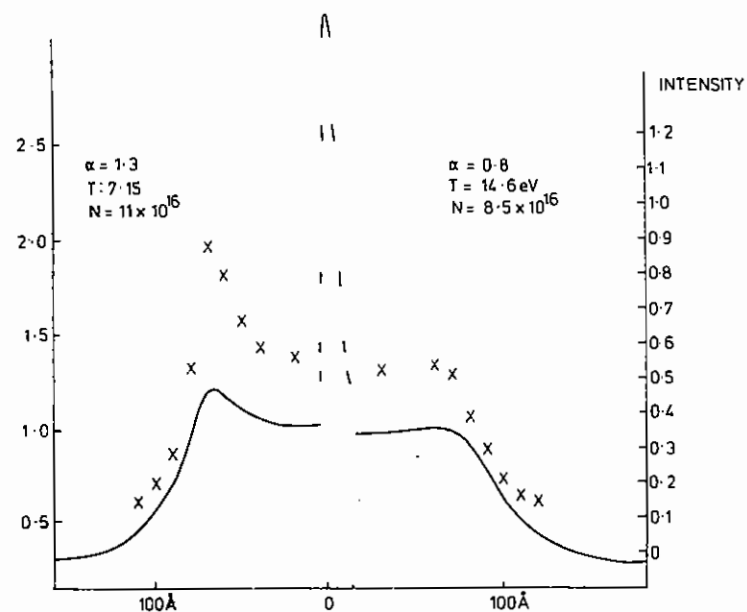
VERTICAL VIEW ALONG Z PINCH AXIS OF INTERACTION REGION

Fig. 5.33 Schematic diagram showing positions of heated and diagnosed volumes

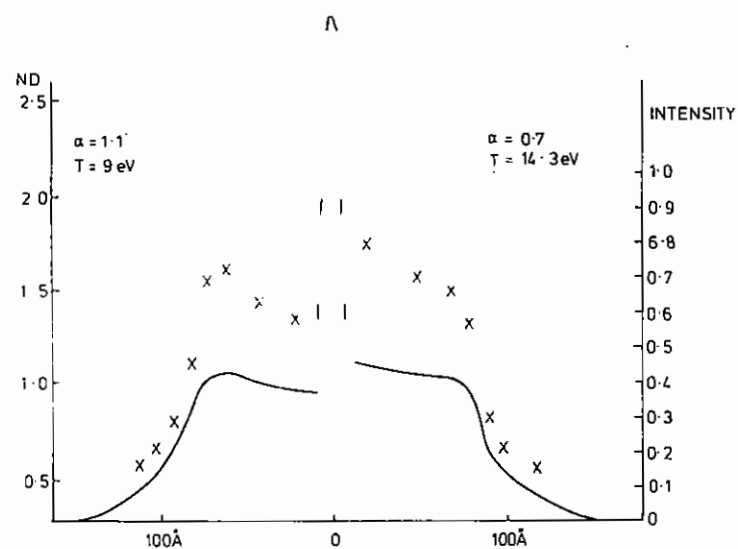
- a) radial view
- b) vertical view.



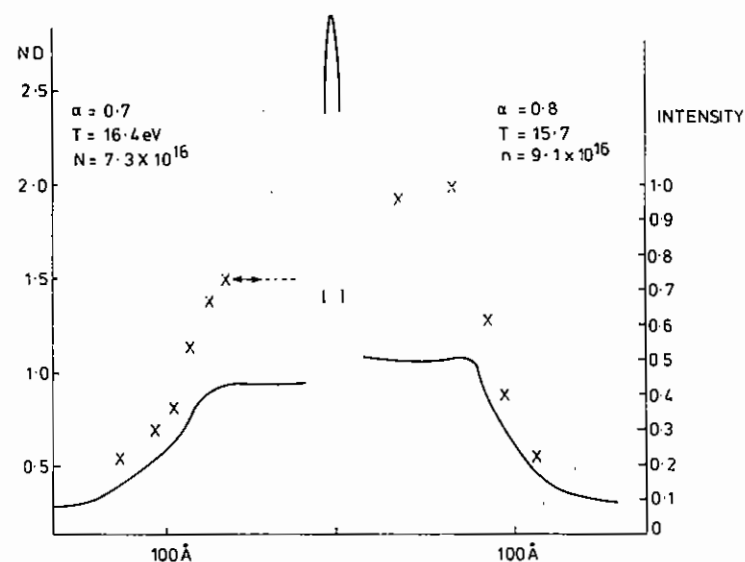
(a)



(b)



(c)



(d)

Fig. 5.34 a) A typical streak photograph is shown. This is microdensitometered perpendicular to the time axis. Microdensitometer traces are shown at three time intervals (solid lines). These are deconvolved for film response (crosses). The profiles each side of the laser wavelength are then fitted to x curves by the method of Kunze. Results are shown on the figures for α , N and T for each side of the spectrum.
b) at 0.5 to 1.0 ns
c) at 1.5 to 2.0 ns
d) at 2.5 to 3.0 ns

asymmetry is evident at all positions with respect to the centre of the heated region. In several instances the asymmetry is seen to relax as V_T decreases. This observation makes it unlikely that the effect could be due to distortion or variations in spectral sensitivity of the camera system.

The reproducibility of the spectra for nominally identical conditions was rather poor, showing particularly noticeable variations in the temperature. The reason for this irreproducibility is not clear. It could be due to varying initial conditions of the z-pinch or due to the relative wandering of the beams.

Interpretation

The measured values of density and particularly the temperature were compared with a one-D cylindrically symmetric MHD code developed at Imperial College by M S White, D Grey and E Wyndham. This code also allows for the non-linear part of the inverse-Bremsstrahlung coefficient of absorption at high irradiances. The temperatures which are observed may be fitted assuming Spitzer conductivity limited at 5% of the free-streaming limit. It is interesting to see that this estimate compares closely with that of D Grey (5.51) where anomalous enhancement of the ion feature was seen and where T_e/T_i was estimated to be approximately four.

The maximum values of λ_e/L inferred from the simulation are about unity, about twice as high as Grey inferred from his observations.

It should be appreciated that this was a trial experiment to establish that the 1 μ m laser beam available at the Rutherford Central Laser Facility was sufficient to generate large values of V_T in a very underdense plasma. The CLF is particularly suitable because of the high quality of the irradiating beams and because of the flexibility of the system. It was also the first time that laser scattering data has been obtained using a streak camera, giving good time-resolved measurements.

A major objective of these experiments is to study the ion feature in

detail and hence obtain information about the effect of ion-acoustic turbulence on the that flow. Information regarding the growth time, decay time and saturation levels as well as the energy spectrum of the turbulence can be obtained from this experiment.

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CHAPTER 6 ABLATIVE ACCELERATION AND COMPRESSION STUDIES

INDEX

- 6.1 INTRODUCTION page 6.1
- 6.2 PULSED X-RADIOGRAPHY OF $\lambda = 0.53 \mu\text{m}$ DRIVEN ABLATIVE IMPLOSIONS page 6.1
 - 6.2.1 Introduction
 - 6.2.2 Experimental Set-up
 - 6.2.3 Data Reduction
 - 6.2.4 Experimental Results
 - 6.2.5 Hydrodynamic Simulations
 - 6.2.6 Conclusions
 - 6.2.7 Ionisation Model of CH Plasma
 - 6.2.8 Self Similar Expansion for a Homogeneous, Dense, Spherical Plasma
- 6.3 STABILITY STUDIES ON ABLATIVELY ACCELERATED RIPPLED FOILS page 6.11
 - 6.3.1 Introduction
 - 6.3.2 Observation of R-T Growth
- 6.4 ANALYSIS OF $\lambda = 0.53 \mu\text{m}$ DRIVEN HOT ABLATIVE IMPLOSION EXPERIMENTS page 6.13
 - 6.4.1 Introduction
 - 6.4.2 Experimental Results
- 6.5 CORE SPECTROSCOPY OF HIGH INTENSITY GREEN IMPLOSIONS AT KMS FUSION page 6.18
 - 6.5.1 Introduction
 - 6.5.2 Experiment and Diagnostics
 - 6.5.3 Results

REFERENCES page 6.22

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6.1 Introduction

Compression results reported in this chapter cover a range of experiments from dominantly exploding pusher (Section 6.5), through implosions with a moderate level of radiative preheat (Section 6.4) to implosions with minimal preheat of the purely ablative type (Section 6.2). These results, together with additional data on Rayleigh-Taylor instability growth rates from rippled plane targets (Section 6.3), have given further valuable information on the parameters necessary for carrying out optimized ablative implosions in terms of laser wavelength and irradiation intensity for good thermal smoothing, target wall thickness requirements to reduce preheating to a tolerable level and shell aspect ratios required to minimise hydrodynamic instability growth. The experiments are a vital prelude to a lengthy campaign of ablative implosions to be carried out in the Spring of 1983 using the enhanced output (500J in 1ns at $\lambda = 0.53\mu\text{m}$) from the recently upgraded laser system (Section 1.3.1):

6.2 Pulsed X-radiography of $\lambda = 0.53\mu\text{m}$ driven ablative implosions6.2.1 Introduction

The efficient implosion of a spherical shell target using laser beam irradiation requires an understanding of the importance of a wide range of parameters and effects. For example appropriate parameters include target geometry ($r, \Delta r$, structure and materials), irradiation conditions (E_L , τ_L , λ , uniformity, focusing) and how these combine to influence effects such as thermal smoothing, preheating and Rayleigh Taylor instability. The experiment described here specifically sets out to study the quality of implosions achieved from the irradiation of low aspect ratio, empty, plastic shell targets with green light ($\lambda = 0.53\mu\text{m}$) at incident irradiance levels of $I \approx 10^{14} \text{ W cm}^{-2}$. We conclude that these targets perform significantly better than in previous experiments reported in earlier annual reports.

6.2.2 Experimental Set-up

Six second harmonic beams ($\lambda = 0.53\mu\text{m}$) delivered typically up to 50Joules of energy onto the target in a pulse of duration 1.3 ns fwhm.

Synchronised to these target drive pulses a seventh fundamental beam ($\lambda = 1.06 \mu\text{m}$) delivered 20-30 J in < 200 psec to a copper backlighting target. The shells were irradiated under focusing conditions such that the main f/1 lenses were focused 3.5 target radii beyond the target centre. Equivalent plane beam images showed the random intensity variation at the target surface to be typically $\times 2$ on a $\approx 20 \mu\text{m}$ scale length. Energy monitors showed the rms deviation between the six incident beam energies to be less than 10% of the average beam energy.

The X-ray emission from the backlighting target was imaged by a $5 \mu\text{m}$ diameter pinhole (in a platinum substrate) onto the CsI photocathode of an X-ray streak camera used in focus mode. The imploding shell was located on the line of sight at a distance $\approx 500 \mu\text{m}$ from the backlighting source and ≈ 1.4 cm from the pinhole (Fig. 6.01). A standard pinhole camera ($\times 4.7$) with two gold pinholes of $5 \mu\text{m}$ and $30 \mu\text{m}$ diameter filtered by 25 μm Be and 25 μm Be plus $5 \mu\text{m}$ mylar respectively was used to view the X-ray emission from each target separately. This gave an indication of the relative brightness of the self emission from the shell target and the uniformity of heating.

The streak camera imaging channel was filtered with 25 μm mylar and 15 μm Be in the cathode substrate giving an average X-ray probe energy of 2 keV with a half intensity energy spread of ≈ 1.0 keV. This is well beyond the lowest energy resonant line emission from neon-like copper at ≈ 1.0 keV and in the analysis described later the backlighting energy spectrum was treated as a continuum.

Irradiated targets were made from polystyrene and had typical diameters in the range $\phi = 120 - 150 \mu\text{m}$, wall thicknesses 9-13 μm and aspect ratios 5-7.5 (Table 6.01).

TABLE 6.01

Shot no., Date	Experimental Parameters						Medusa			• Backlighting			Detection		
	ϕ μm	ΔR μm	$r/\Delta r$	$E_{2\omega\text{inc}}$ on target J	$I_{2\omega\text{inc}}$ $\times 10^{13}\text{Wcm}^{-2}$	D/R	I_{inc} $\times 10^{13}\text{Wcm}^{-2}$	$E_{\text{abs.}}$ J	$I_{\text{abs. (Peak)}}$ $\times 10^{13}\text{Wcm}^{-2}$	ϕ μm	E_{ω} J	I_{inc} $\times 10^{13}\text{Wcm}^{-2}$	Filter, thickness	Gain stop w.r.t. Max.	Delay w.r.t. Peak of the Pulse ns
05,26/8/82	139	11.2	6.2	32.5	4.24	3.8	3.5	20.2	2.9	350	31.4	8.55	mylar 25 μ Be 15 μ	-2	0
08,26/8/82	132	12.6	5.2	32.5	4.56	3.8	4.8	21.5	10.3	350	26.5	7.22	mylar 25 μ Be 15 μ	-2	0
10,26/8/82	132	9.2	7.2	40.0	5.63	3.8	4.6	20.4	14.2	350	24.1	6.56	mylar 25 μ Be 15 μ	-2	0.2
17,26/8/82	149	11.9	6.3	34.4	3.79	3.8	3.1	22.0	2.7	350	24.4	6.65	mylar 25 μ Be 15 μ	-2	0.2
21,26/8/82	126.65	8.5	7.5	39.4	6.01	3.8	4.9	20.7	7.0	350	22.5	6.13	mylar 25 μ Be 15 μ	-2	0.4
01,27/8/82	118.51	9.2	6.4	44.8	7.81	3.8	6.6	18.4	30.5	350	34.7	9.45	mylar 25 μ Be 15 μ	-2	0.4
03,27/8/82	123.69	9.4	6.6	42.4	6.79	3.8	5.8	19.6	28.8	350	26.1	7.11	mylar 25 μ Be 15 μ	-2	0.6
06,27/8/82	118.51	9.1	6.4	36.8	6.42	3.8	5.4	16.9	29.8	350	31.1	8.47	mylar 25 μ Be 15 μ	-2	0.7
08,27/8/82	144.43	10.1	7.2	30.9	3.62	3.5	3.1	20.1	8.7	350	23.6	6.43	mylar 25 μ Be 15 μ	-2	0.7
07,02/9/82	145.4	11.3	6.5	40.5	4.56	4.25	3.8	24.0	3.2	350	20.7	5.64	mylar 25 μ Be 15 μ	-2	-0.3
10,02/9/82	124.43	10.6	5.9	44.2	6.99	3.8	5.7	22.4	26.8	350	24.8	6.75	mylar 25 μ Be 15 μ	-2	-0.5
15,02/9/82	134	11.9	5.6	43.1	5.87	3.8	4.8	22.0	14.9	350	23	6.26	mylar 25 μ Be 15 μ	-2	1.3
17,02/9/82	143.7	11.13	6.5	43.9	5.21	3.8	4.4	25.3	12.0	350	22	5.99	mylar 25 μ Be 15 μ	-2	1.6
19,02/9/82	151.1	11.02	6.9	52.0	5.57	3.8	4.6	28.8	12.9	350	28.5	7.76	mylar 25 μ Be 15 μ	-2	1.8

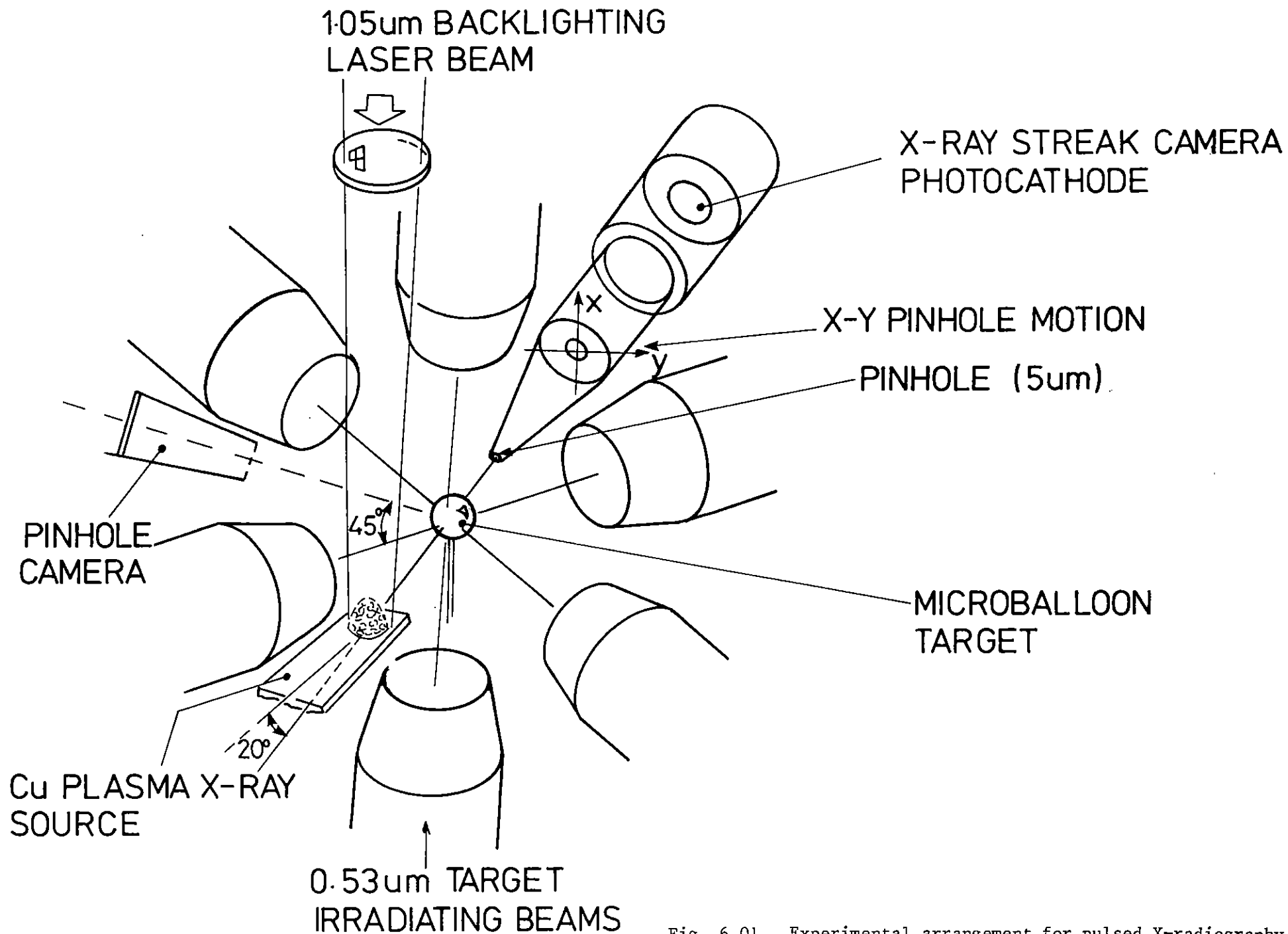


Fig. 6.01 Experimental arrangement for pulsed X-radiography.

6.2.3 Data Reduction

2-D images showing absorption features were recorded on HP5 film and processed to transform the information recorded into a mass density distribution in the target plane. Each original HP5 film record was microdensitometered and an extrapolation made to simulate the backlighting plasma image density in the absence of any absorption. Using a calibration curve for HP5 film the transmission of the target to backlighting X-rays was plotted as a function of radius. This profile was divided into typically 10 radial elements and input into an Abel inversion program (POLYABEL) which took typically 4 minutes to run on an Apple microcomputer. As input data POLYABEL requires the experimental conditions r , ΔR , I_{abs} , λ , τ_L and the time into the event when the pulsed 2-D image was recorded (τ_{probe}). It also requires a single temperature T_{BL} which characterises an exponentially decaying distribution of the backlighting source spectrum. i.e. $I(E) = I_0 \exp(-E/T_{\text{BL}})$ where E is the X-ray photon energy. The target material is specified in terms of its initial density, ρ_0 , and a mass absorption coefficient, μ_1 , corresponding to $E = 1$ KeV. The program takes $\mu(E) = \mu_1/E^3$. The X-ray filter in front of the imaging system is specified in terms of its thickness, t_a , density, ρ_a , and mass absorption coefficient, μ_{1a} . The transmission through a target with a radial density profile $\rho(r)$ is then given by

$$T(r) = \frac{\int \exp\{-E/T\} \exp\{-(\mu_{1a}\rho_a t_a + 2\mu_1 \int_0^R f(R) dR)/E^3\} dE}{\int \exp\{-E/T\} \exp\{-(\mu_{1a}\rho_a t_a)/E^3\} dE} \quad (6.01)$$

where $f(R) = \frac{\rho(R)R}{(R^2 - r^2)^{1/2}}$, R_0 is the boundary radius of the absorption feature (i.e. $T(R_0) = 1.0$) and the limits of the energy integration are typically 0.5-5 keV, with the peak of the probe spectrum at $E \approx 2$ keV. A subroutine in POLYABEL calculates transmission for a suitable range of

density and length products in the form of an array and in the inversion procedure equation 6.01 is solved by interpolation of the array elements to find $\int_0^{R_0} f(R)$ for an experimental value of $T(r)$. Working inwards from the outer shell the density of each shell can be deduced. (Although all the data here deals with empty pure polymer shells the programme can handle simple composite targets e.g. gas filled single material shells or double layered empty shells.)

POLYABEL calculates the residual target mass at the probing time through

$$M_a = 4\pi r^2 \left\{ \Delta r \rho - \int_0^t \dot{m}(t) dt \right\}$$

where $\dot{m}(t) = 0.85 \times 10^5 I_a b_s^{0.72}(t)$ is the ablation rate at $\lambda = 0.53 \mu\text{m}$ and the laser pulse profile is taken as triangular. This mass M_r , is compared to the mass deduced from the Abel inversion density profile

$$M_a = 4\pi \int_0^{R_0} \rho(r) r^2 dr$$

which depends on the input parameters T_{BL} and μ_1 and the agreement of the fit is expressed as $\epsilon = (M_a(T, \mu_1) - M_r)/M_r \times 100\%$. Hence acceptable density profiles can be limited to a region of T_{BL} and μ_1 parameter space, which itself has other restrictions as discussed later.

6.2.4 Experimental Results

Time integrated pinhole images of the backlighting source and the shell self emission showed that at a filtered energy of ≈ 1.4 keV the shell emission intensity was $< 20\%$ of the source intensity. The source was viewed at $\approx 30^\circ$ to the backlighting substrate normal and the contrast should be improved as viewed by the streak camera i.e. almost edge-on into the plasma. An attempt to record the shell self-emission only with the X-ray streak camera pinhole showed the CH self emission to be undetectable with 2 keV filtering and thus no correction was used in the analysis.

Figure 6.02 shows a radiograph of an unirradiated shell and provides the overall magnification and resolution of the imaging system (i.e. $\times 16$ and

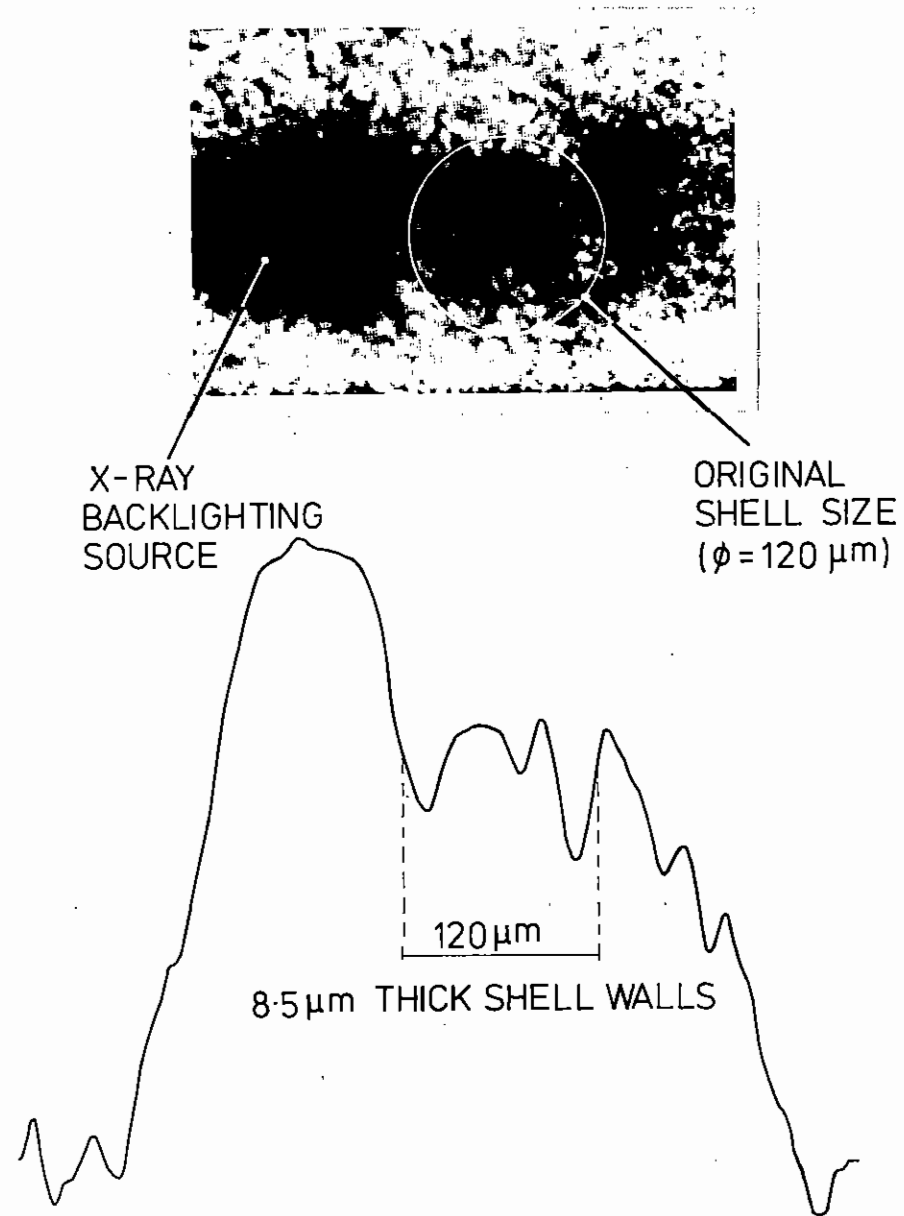


Fig. 6.02 X-radiograph and densitometer trace of unirradiated plastic shell target.

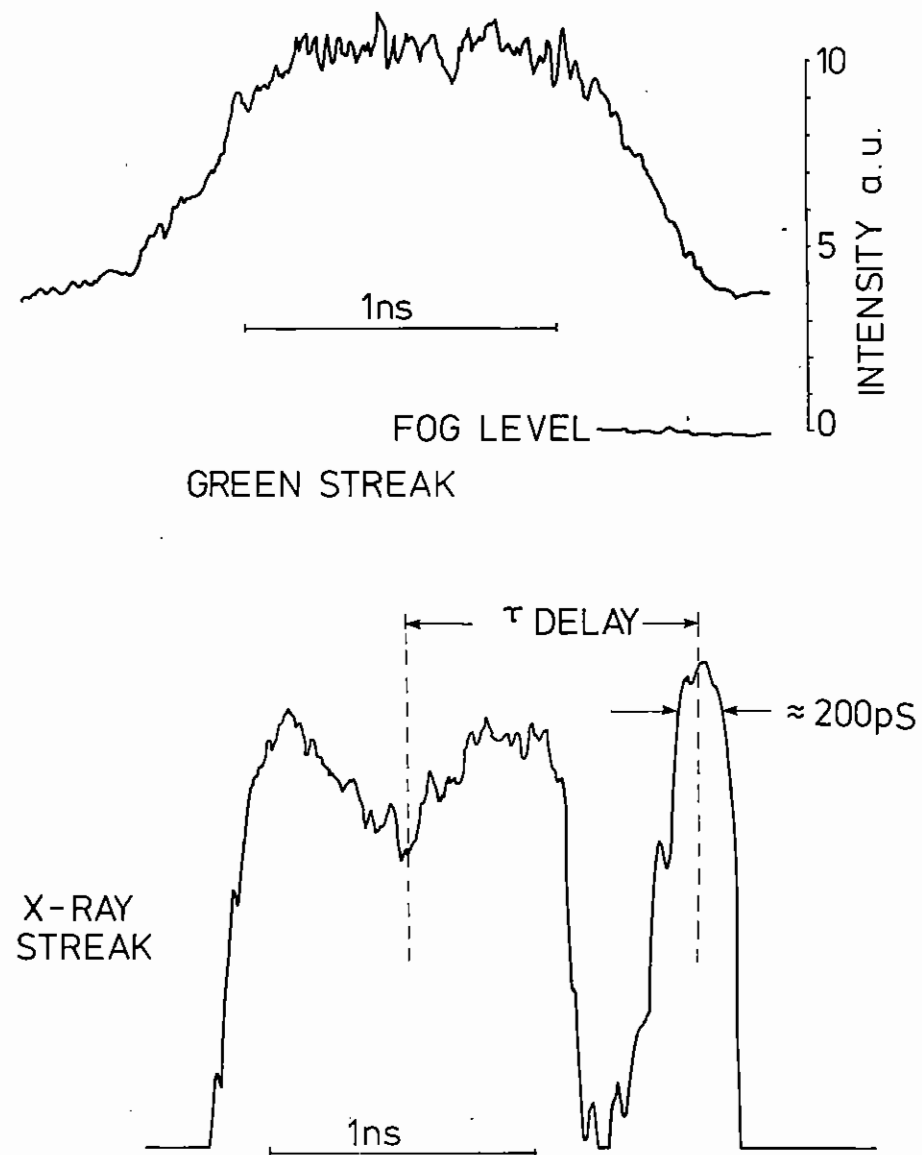


Fig. 6.03 Densitometer trace of a streak camera record of the X-ray emission from a copper target irradiated with the 1ns driver beam and the short pulse backlighting beam.

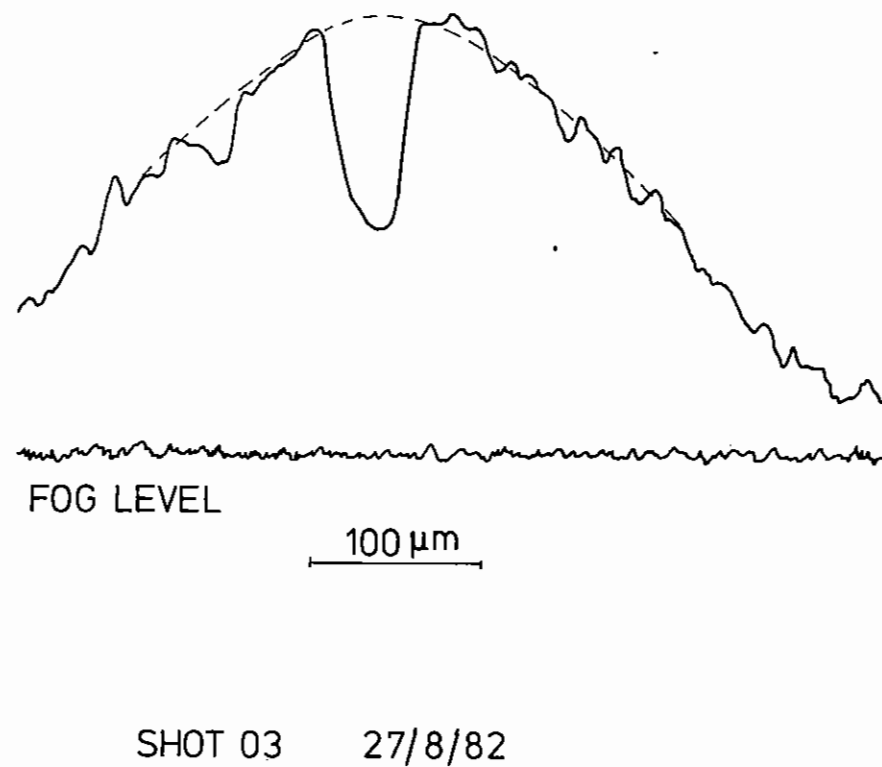
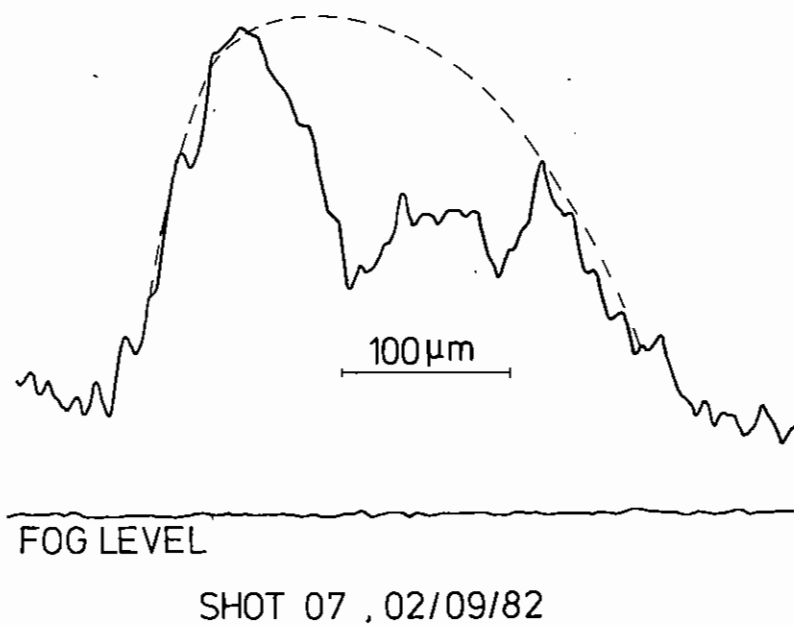
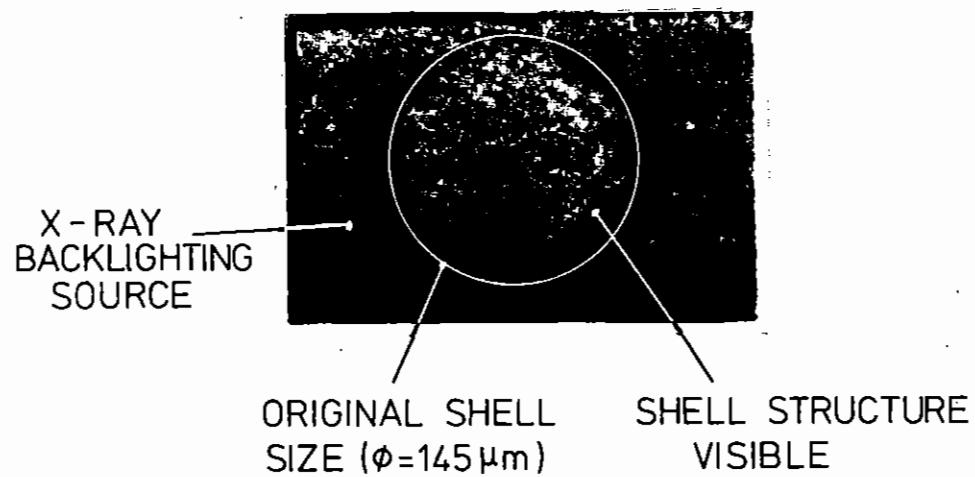


Fig. 6.05 Densitometer trace of X-radiograph of imploding shell target takes class to peak compression.

Fig. 6.04 Pulsed X-radiograph of imploding shell target taken early in the implosion.

10 μm in the target plane respectively.)

Figure 6.03 shows a very saturated X-ray streak of the emission from a flat copper target irradiated with the normal backlighting beam and one of the six shell driver beams. Filtering is the same as for radiography. This streak gives the fwhm duration of the X-ray probe beam (< 200 psec) and gives a calibration of the delay between the probe pulse and the peak of the shell driver pulse assuming the X-ray emission is symmetric about the time of peak laser intensity. (This particular shot showed temporal structure in the long pulse due to mode-beating.)

Figure 6.04 shows a radiograph taken early in the implosion and illustrates that the shell is remaining intact. X-ray transmission measurements through the centre of a shell image can be readily deconvolved to give an estimate of the source temperature (T_{BL}) as the absorbing material is effectively planar and two shell walls thick. Four estimates of this type were possible indicating a single temperature source characteristic of $T_{BL} = 460 + 40 \text{ eV}$ with an incident intensity scaling trend of the order $T_{BL} \propto I^{0.13}$.

Figure 6.05 illustrates a radiograph taken close to the peak compression phase. The U-shape of the microdensitometered absorption feature (as opposed to V-shaped seen in earlier experiments) indicates that the core material is contained in a well defined volume of space.

Results from Abel inversion fitting of the radiographs, described earlier using POLYABEL, are illustrated in Figures 6.06, 6.08 and 6.10. These correspond to radiographs taken at probe times $\tau_p < \tau_{imp}$, $\tau_p \approx \tau_{imp}$ and $\tau_p > \tau_{imp}$ where τ_{imp} is the time of peak compression predicted by Medusa for a given shot. The shaded regions correspond to fits where the Abel integrated core mass and the estimated shell residual mass agree to within 20% for assumed source temperatures in the range $T_{BL} = 350 - 550 \text{ eV}$. Figures 6.07, 6.09 and 6.11 show the corresponding range of Abel inverted core density profiles. The effective mass absorption coefficient at 1 keV (μ_1) needed for the shaded fit regions consistently showed that at probe times corresponding to peak

compression (μ_1) eff $\ll$ (μ_1) for cold material while fits for radiographs taken well after peak compression required (μ_1) eff $<$ (μ_1) for cold material.

6.2.5 Hydrodynamic Simulations

Medusa hydrodynamic simulations were run for each shot which provided a useful radiograph. The following procedure was used to make comparisons with experiment. A Medusa simulation gave times of peak compression T_{IMP} and peak laser intensity T_L . The experimental probe times were measured as delays T_D with respect to the peak of the laser pulse generated X-ray emission from a flat target. (T_D is $\pm$ depending on whether probing is late or early with respect to laser pulse peak.)

The Medusa simulation time used for comparison with a given shot was then taken as $T_M = T_L + T_D$ and T_M/T_{IMP} was used as a normalising time in plotting data to indicate when a radiograph was taken with respect to the peak compression time.

The density length product ($A_{\frac{1}{2}}$) through a chord of target material required to give a 50% transmission of the X-rays from a backlighting source of temperature T_{BL} with the experimental filtering can be estimated from equation 6.01 using cold material mass absorption coefficients and is plotted in Fig. 6.12. Medusa $\rho(r, t)$ outputs were post-processed to find the chord giving an integrated density length product equal to $A_{\frac{1}{2}}(T_{BL})$. Hence the trajectory of the 50% transmission chord ($r_{\frac{1}{2}}(t)$)_{MED} can be found as a function of time and in particular it can be compared to the experimental radius for 50% transmission at $t = T_M$. Figure 6.13 shows a typical normalised $(r_{\frac{1}{2}}/r_0)_{MED}$ trajectory. At early times, when the shell structure remains intact $r_{\frac{1}{2}}$ comparison with Medusa was made at the midshell radius as the experimental spatial resolution was comparable to the shell thickness. Measured radii of the experimental absorption dips due to the shell were corrected slightly to allow for a weighting of the shell absorption profile towards the inside edge.

SHOT 10 26/08/82

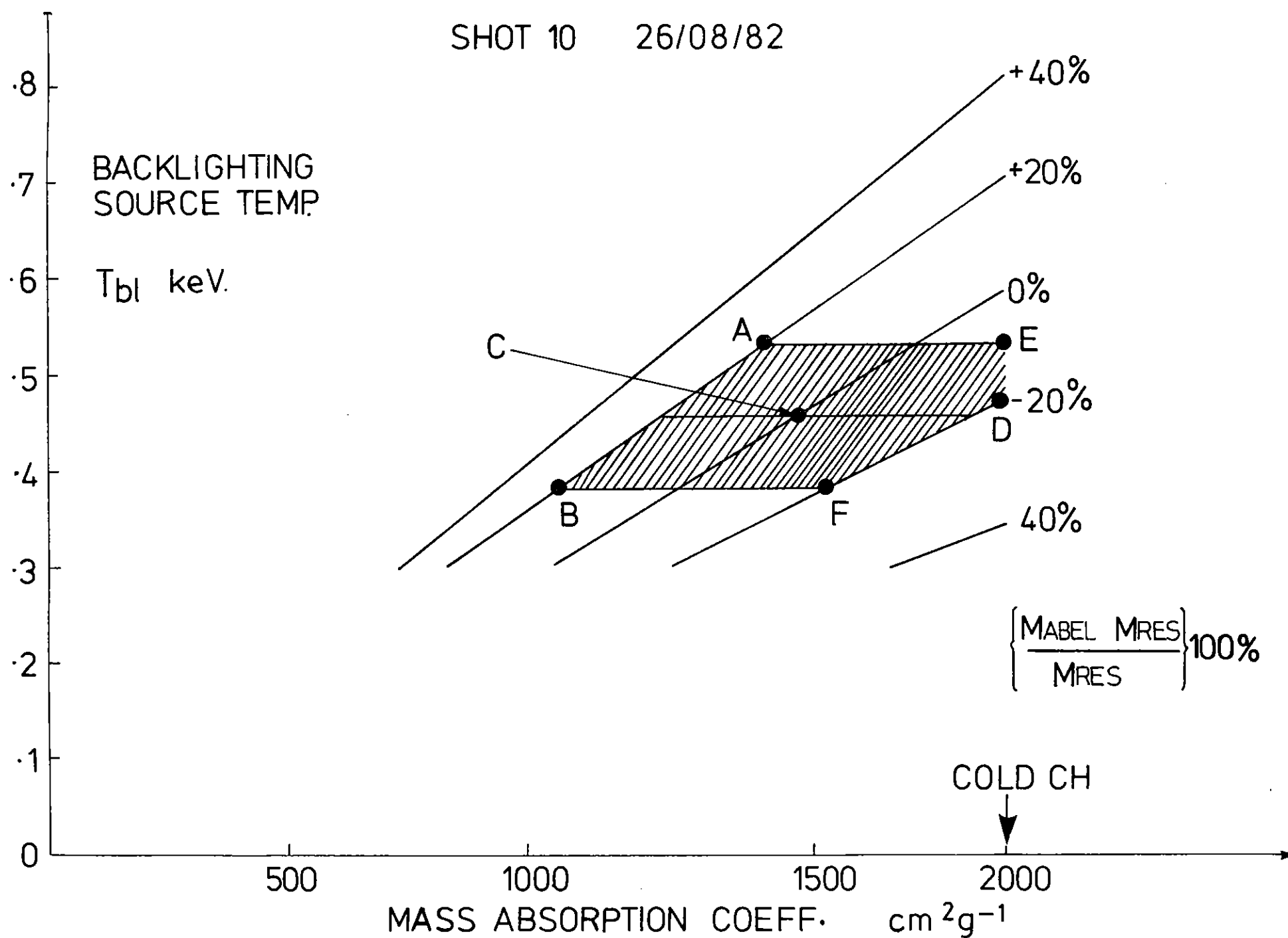


Fig. 6.06 Results from Abel inversion fitting of X-radiograph for probe time $\tau_p < \tau_{imp}$ showing region in which the integrated core mass and the shell residual mass agree to within 20% for source temperatures in the range 350-550 eV.

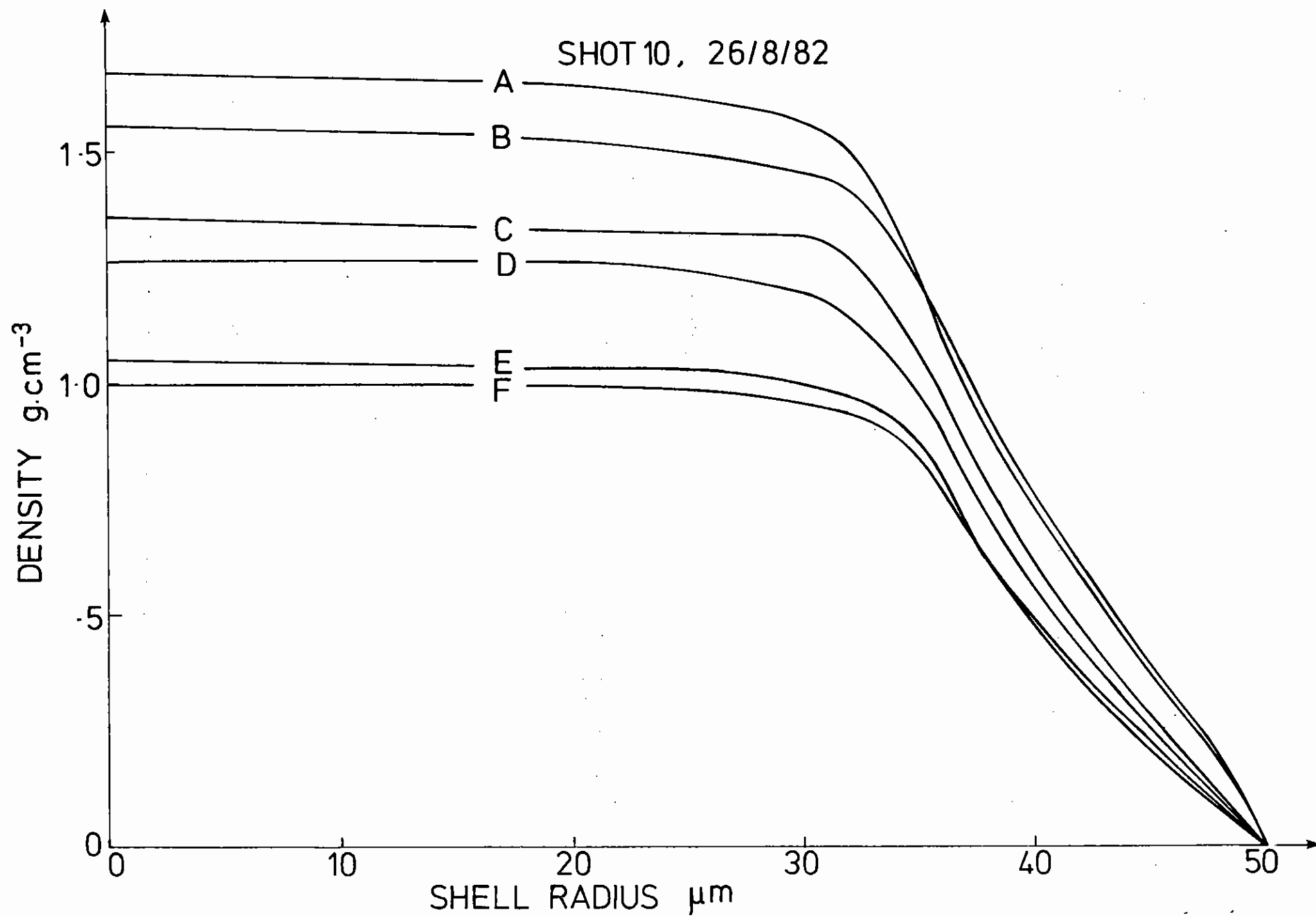


Fig. 6.07 Range of Abel inverted core density profiles corresponding to the points indicated in Fig. 6.06.

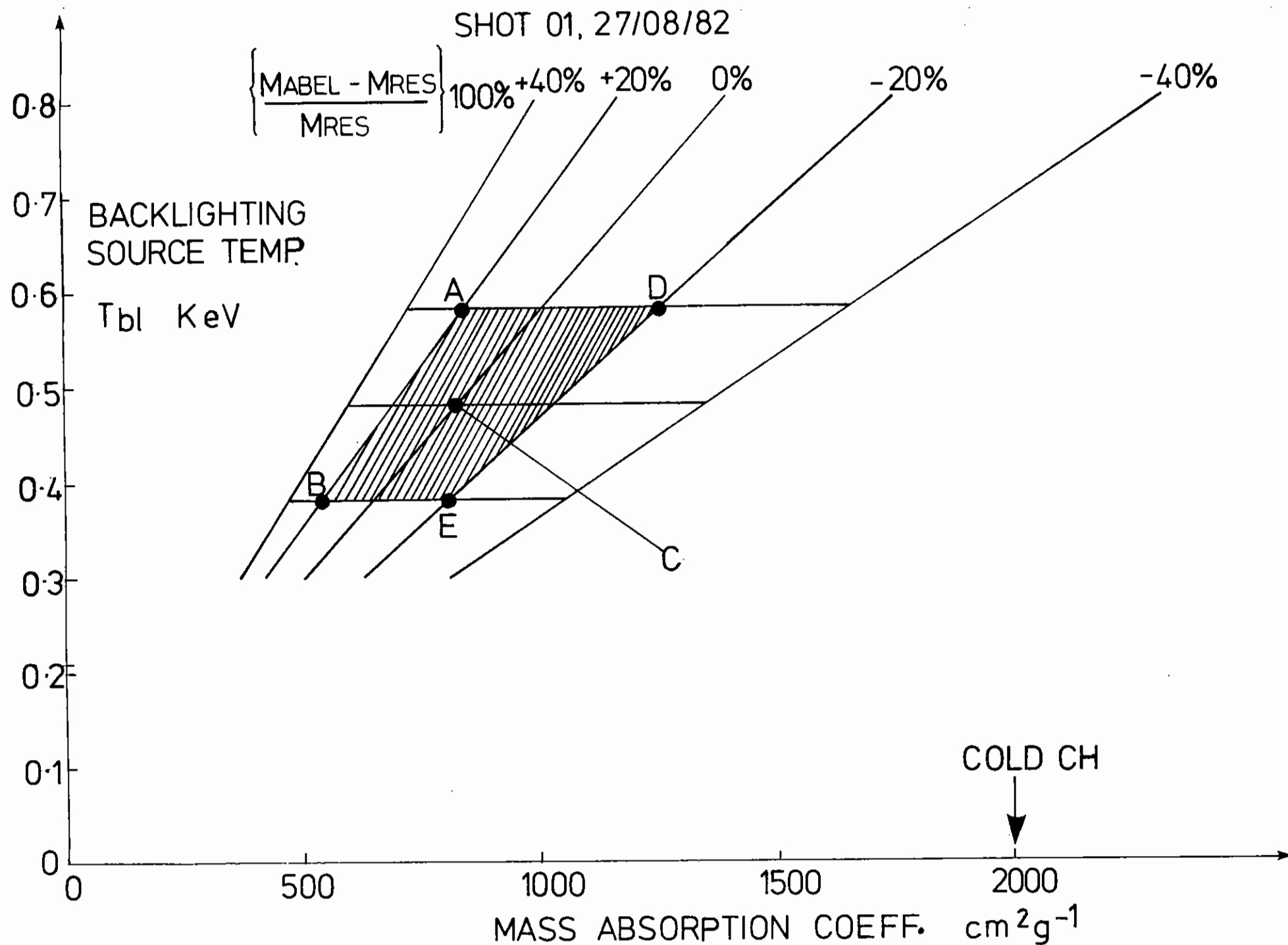


Fig. 6.08 As figure 6.06 only for a probe time close to peak compression ($\tau_p \sim \tau_{imp}$).

SHOT 01 27/8/82

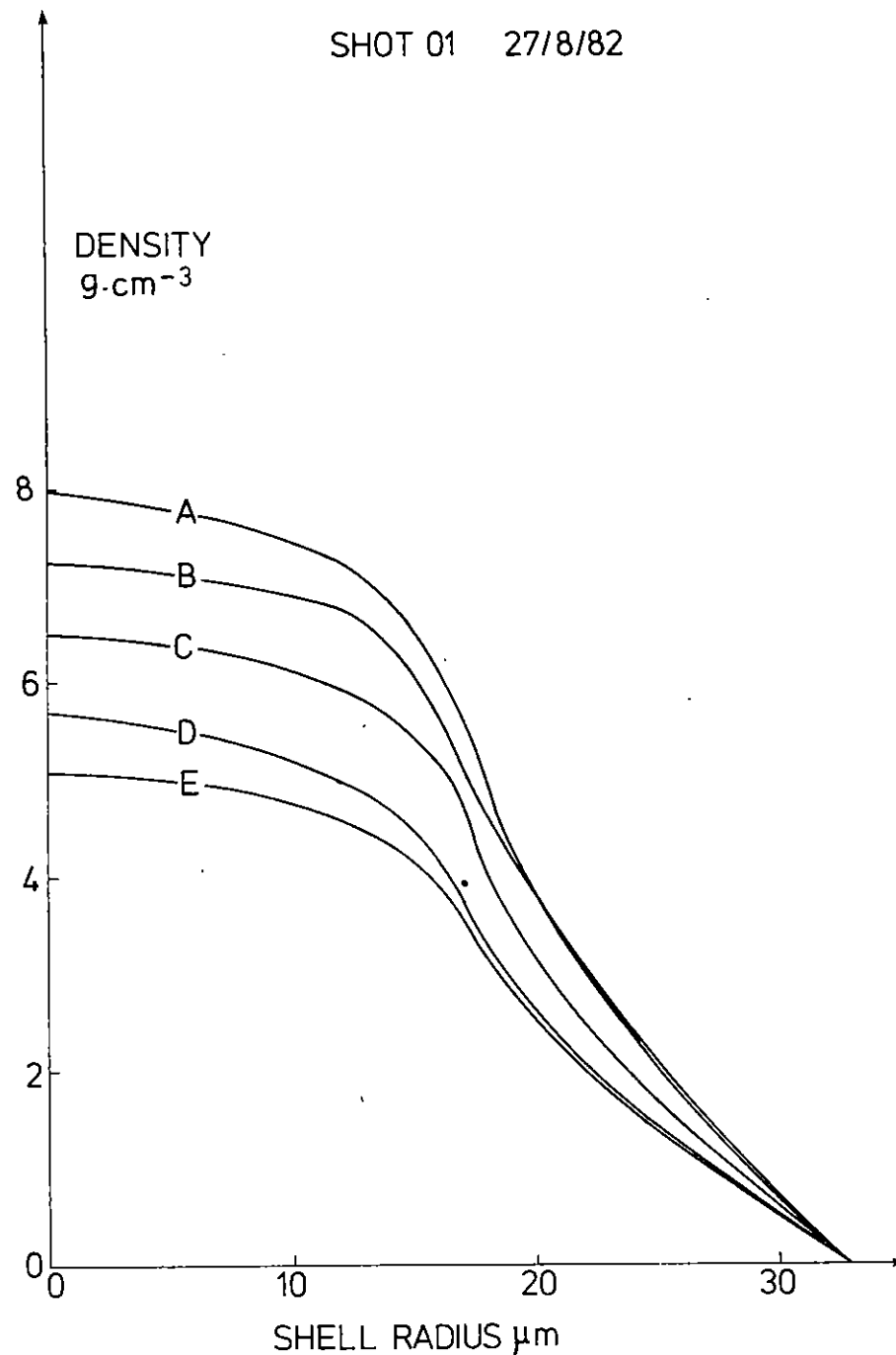


Fig. 6.09 Core density profiles corresponding to the data in Fig. 6.08.

6.2.6 Conclusions

As illustrated in Fig. 6.14 we have found good agreement between $(r_{1/2})_{MED}$ and $(r_{1/2})_{EXP}$ for radiographs taken at times up to and including the peak compression phase. This agreement is significantly better than from similar data taken earlier with thin walled, higher aspect ratio shells and lower irradiances and reported in previous annual reports. The improvement in performance is attributed to longer implosion times and higher irradiance both enhancing thermal smoothing providing a more uniform driving ablation pressure at the target surface. In addition the lower aspect ratio, thicker wall targets reduce the level of radiative preheat at the inner regions of the shell making the volumetric compression at later times greater.

Two sources of information suggest that the temperature of the core material is significantly higher than Medusa predicts (run without radiation transport). These are (1) the low values of $(\mu_1)_{eff.}$ needed to obtain reasonable fits in the Abel inversion procedure and (2) the discrepancy of the $(r_{1/2})_{EXP}$ and $(r_{1/2})_{MED}$ values at times after the peak compression phase. In these two cases we find $(\mu_1)_{eff.}$ can be as low as 40% of the cold material value and $(r_{1/2})_{EXP}$ can be up to 50% larger than the Medusa simulation respectively. Both observations can be attributed to the core at peak compression reaching average temperatures in the region of 150 eV as compared to typically 20 eV in the simulation. A temperature of ≈ 150 eV is broadly consistent with an adiabatic compression of the initial shell material at a preheated temperature of $\lesssim 10$ eV, a value estimated from shock preheating and radiative preheating obtained from absolute X-ray emission measurements made of the continuum flux. The temperature of 150 eV was estimated from two simple models, described in Section 6.2.7 and 6.2.8. These discuss the temperature dependence of the mass absorption coefficient due to ionisation of carbon K-electrons and the time dependence of the absorption profiles expected from adiabatically expanding plasma cores of different initial temperatures.

Table 6.02 summarises the comparison between experiment and simulation.

SHOT 17 02/09/82

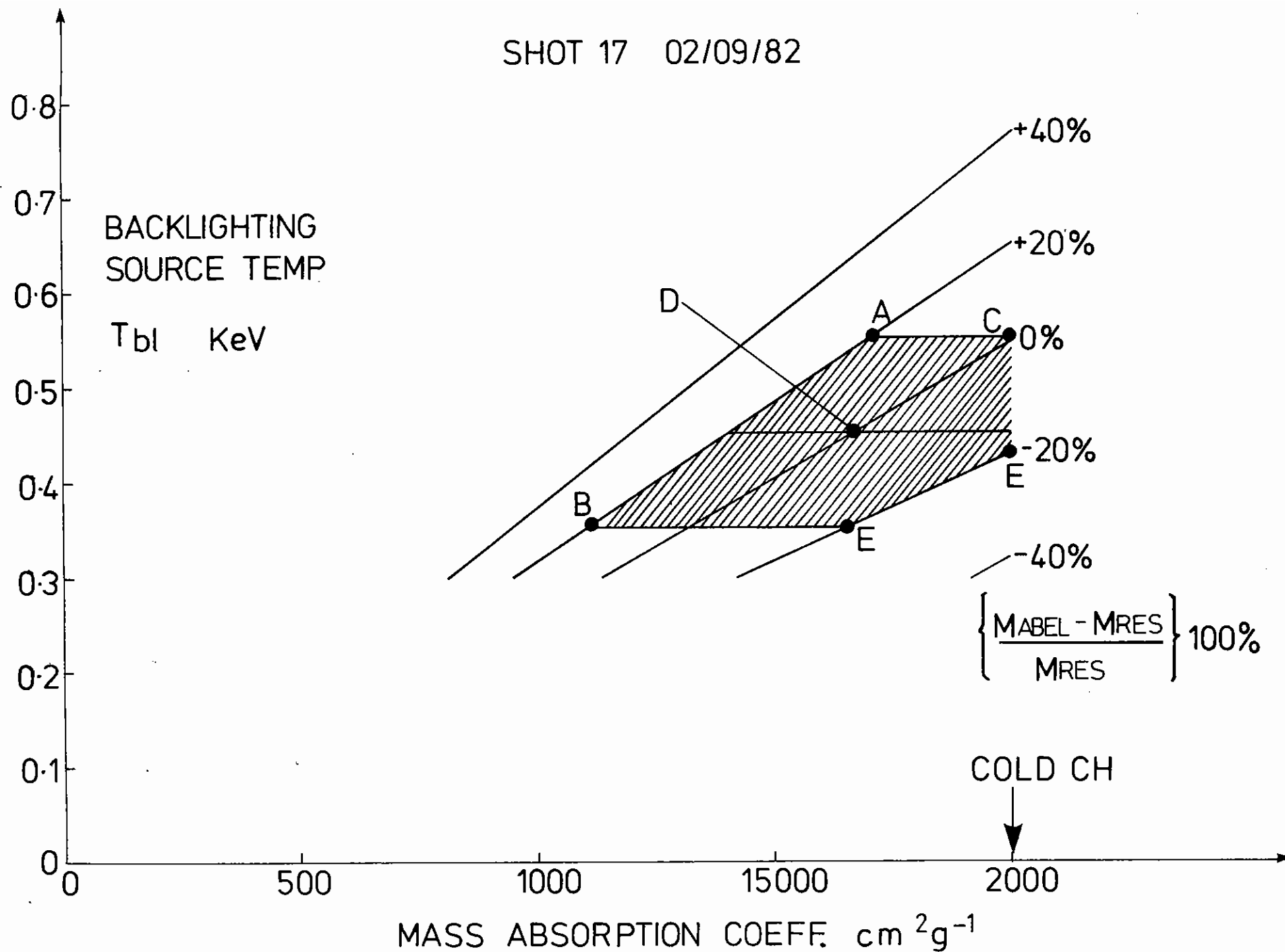


Fig. 6.10 As figure 6.06 only for a probe time after peak compression ($\tau_p > \tau_{imp}$).

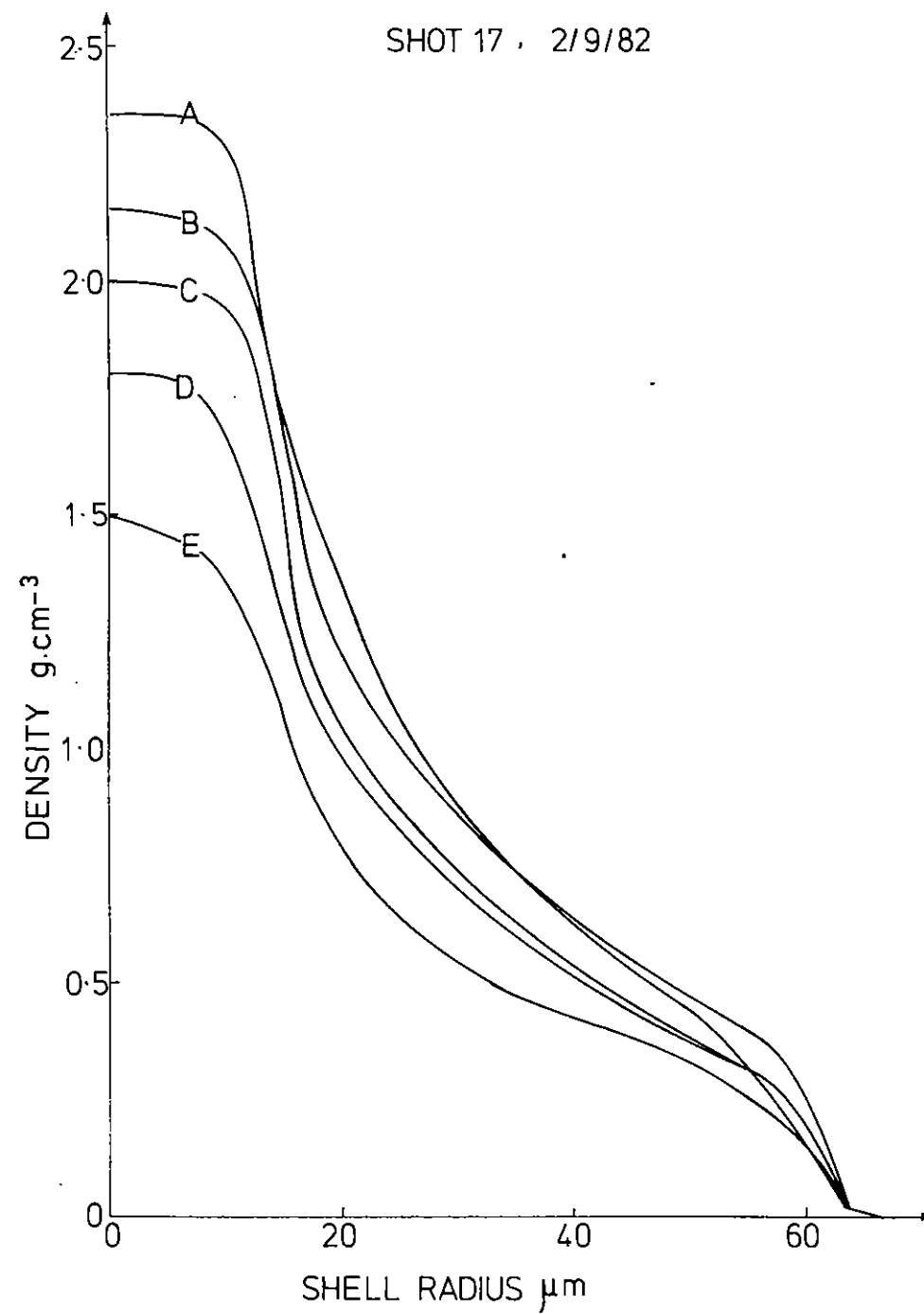


Fig. 6.11 Core density profiles corresponding to the data in Fig. 6.10.

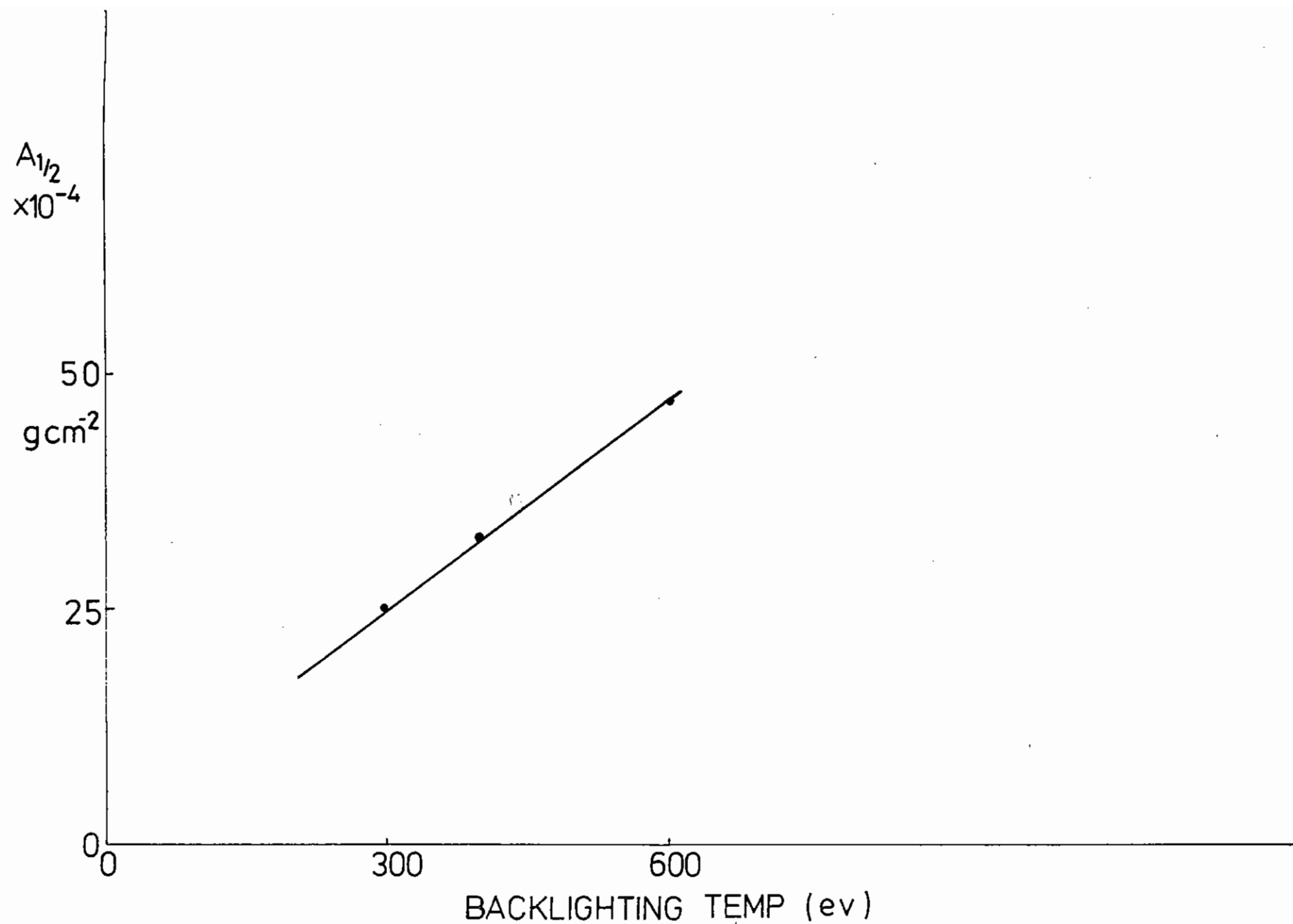


Fig. 6.12 Density x length product ($A_{1/2}$) to give 50% X-ray transmission for various source temperatures calculated from equation 6.01.

The average densities agree well at all times throughout the implosion, where ρ_{av} is defined as

$$\int_0^{r_1} \rho(r) 4\pi r^2 dr / \int_0^{r_1} 4\pi r^2 dr = \rho_{av}$$

The discrepancy in the peak densities at peak implosion time is a result of spatial and temporal smearing of experimental data due to finite resolution. The average Medusa temperature is defined as

$$(T_e)_{av} = \int_0^{r_1} \rho(r) T_e(r) 4\pi r^2 dr / \int_0^{r_1} 4\pi r^2 dr = (T_e)_{av}$$

No attempt was made to include spatial and temporal resolution in the Abel inversion procedure or to iterate effective absorption coefficients as a function of density but this will not effect the main conclusions. We have observed core densities $\rho > 5 \text{ g cm}^{-3}$ and Pr products of $\sim 10^{-2} \text{ g cm}^{-2}$. Finally we can point out that X-ray backlighting at moderate probe energies can be carried out to high densities if the high density material is hot enough to have significant K-shell ionisation. However this would require independent measurements of core temperatures, perhaps through seeded impurities, and more sophisticated modelling of radiation transfer processes in these non-ideal plasmas.

6.2.7 Ionisation model of CH plasma

Saha's equation was used to find the fractional ionisation abundances of carbon ions for a range of plasma densities and temperatures. The hydrogen atoms were assumed ionised at all times. Iteration of the solution to give a self-consistent average charge per atom gave the fraction of each ion species present at a given density and temperature. Higher densities suppress the ionisation level of the plasma at a fixed temperature. From the ionisation fractions and cross-sections for photoionisation of hydrogen and helium-like ions an effective mass absorption coefficient was estimated for 1 keV photons through

$$(\mu_1)_{eff} = n_{oc} \rho (\sigma_1(H) f(H) + \sigma_1(He) f(He) + \sigma_{cold} \sum f(rest))$$

where the cross-sections for H and He like ions were $2.4 \times 10^{-20} \text{ cm}^2$ and $4.0 \times 10^{-20} \text{ cm}^2$ respectively and the other plasma ions were treated as natural absorbers with $\sigma_{cold} = 4.43 \times 10^{-20} \text{ cm}^2$. The fractional abundances are

$$f(z) = \frac{n_c(z)}{n_{oc}} \text{ with } n_{oc} = \frac{\rho}{13m_p} \text{ cm}^{-3}, \text{ where } m_p \text{ is the proton mass}$$

The model predictions for effective mass absorption coefficients is shown in Figure 6.15 and indicates that the absorption coefficients needed for some Abel inversion fits correspond to plasma densities in the region of $\sim 1 \text{ g cm}^{-3}$ at temperatures $\approx 150 \text{ eV}$.

6.2.8 Self similar expansion model for a homogeneous, dense, spherical plasma

Assume the compressed core has uniform density ρ_0 , temperature T_0 and boundary radius r_0 at a time $t = 0$ when conversion of plasma thermal energy to kinetic energy of radially flowing plasma begins. If the expansion is self-similar and adiabatic the plasma remains homogeneous such that at times, t , the plasma has a density ρ , temperature T and a boundary at radius r which is expanding at a velocity v . These variables are related by consideration of plasma energy conservation which yield

$$r(t) = r_0 + \int_0^t v(t) dt \quad (6.02)$$

SHOT 3 27/8/82

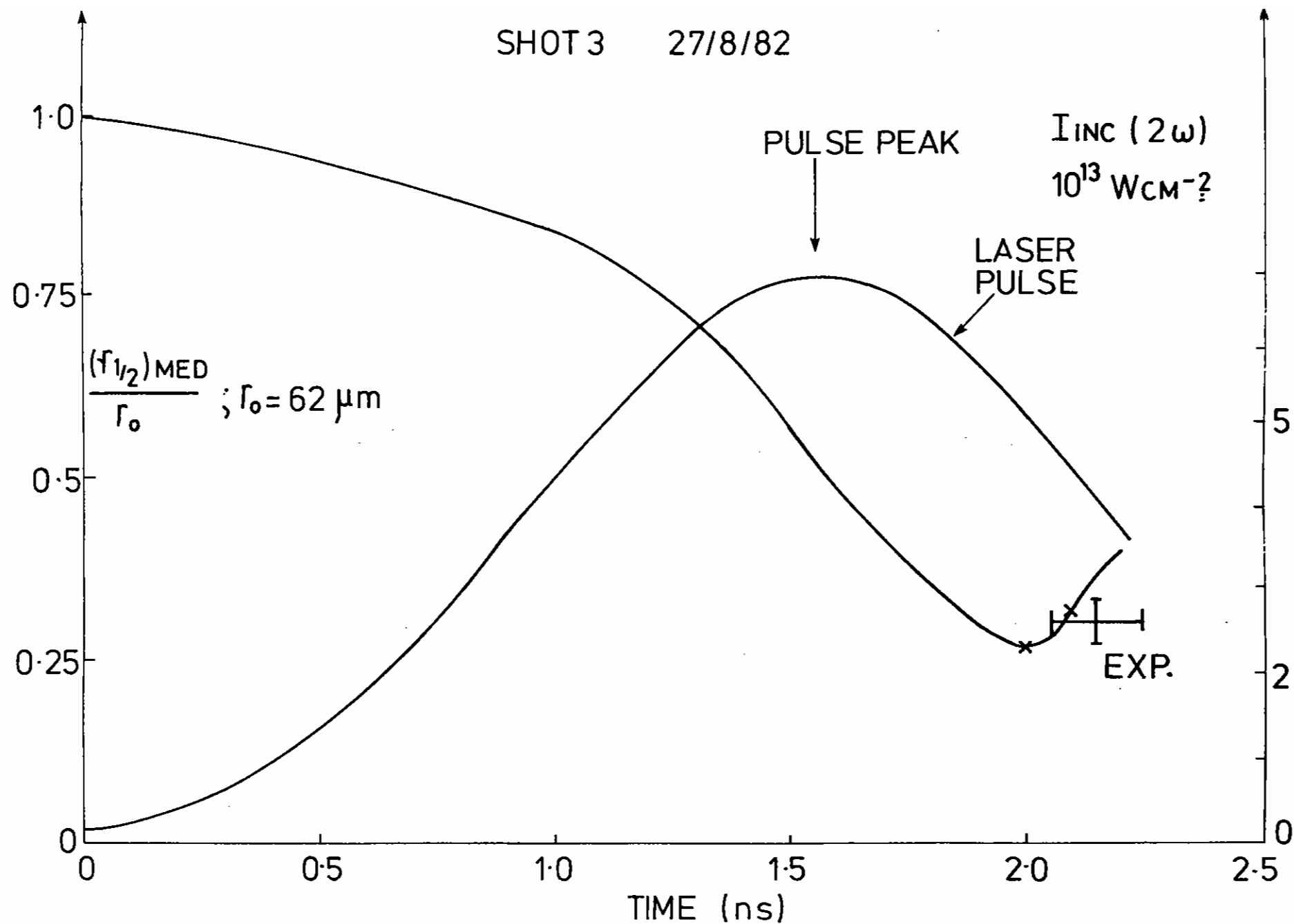


Fig. 6.13 Normalised shell trajectory ($r_{1/2}/r_0$) obtained from MEDUSA for a typical shot showing comparison with experimental observation.

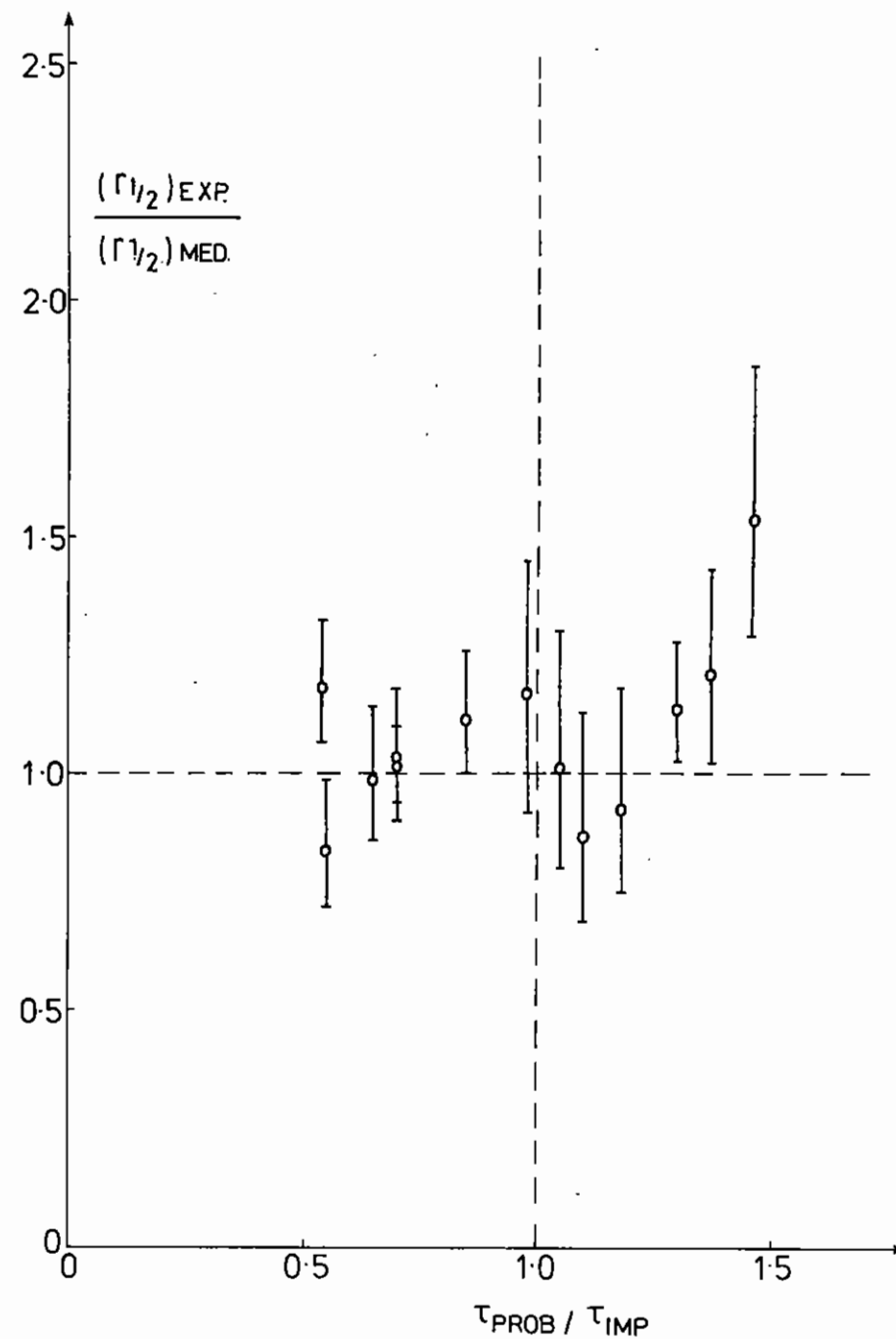


Fig. 6.14 Ratio of experimentally measured and computed core radii at various X-ray probe times before and after peak compression.

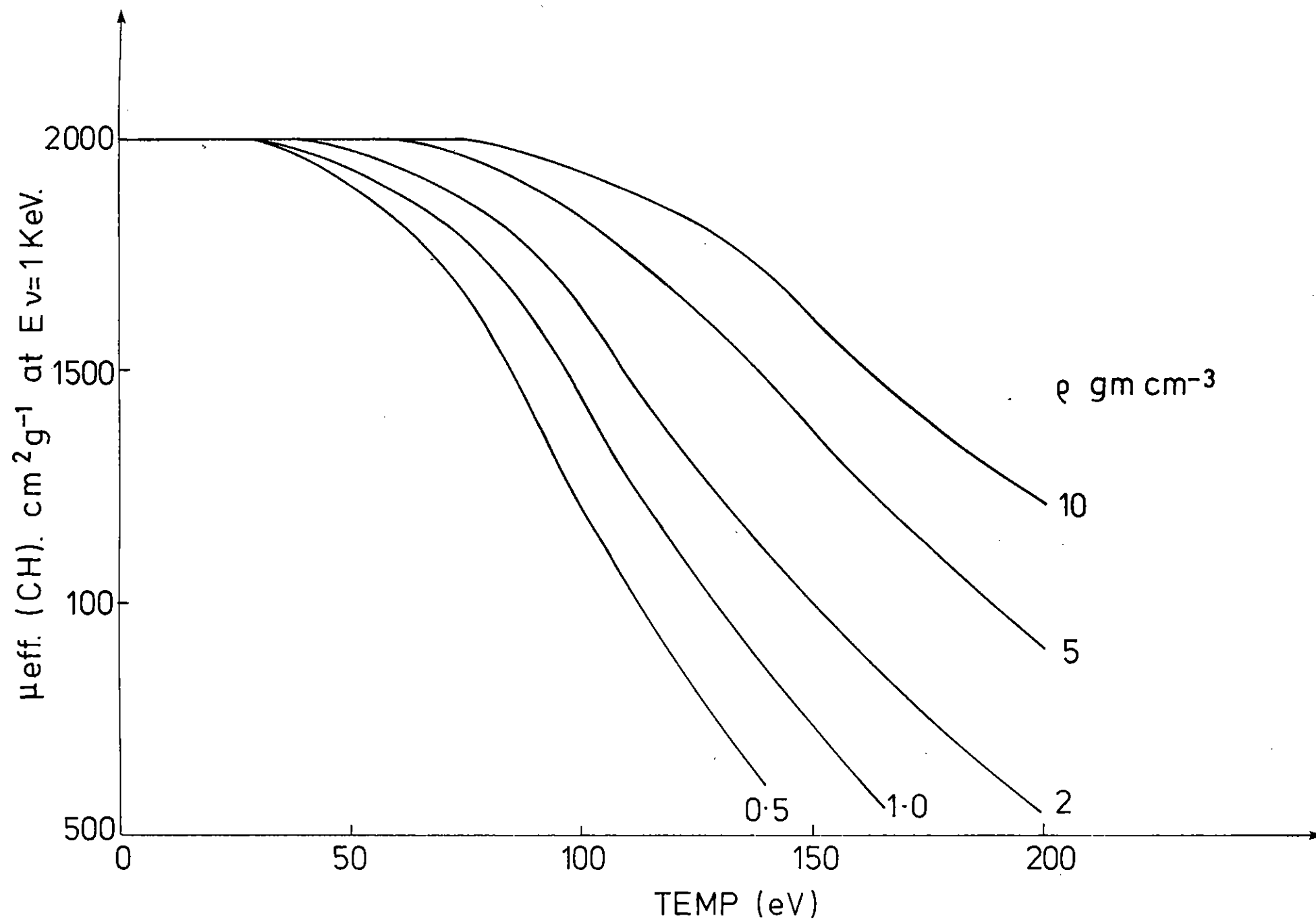


Fig. 6.15 Ionization model predictions for effective absorption coefficients for various core plasma densities and temperatures.

$$v(r) = v_{\infty}(1 - (r_0/r)^2)^{\frac{1}{2}} \quad (6.03)$$

$$\rho = \rho_0 (r_0/r)^3 \text{ and } T = T_0 (r_0/r)^2 \quad (6.04)$$

where v_{∞} is the asymptotic expansion velocity given by $v_{\infty} = \sqrt{CT_0}$ and C is a constant. Combining equations (6.02) and (6.03) solving gives

$$v(t) = \frac{v_{\infty}^2 t}{(r_0^2 + v_{\infty}^2 t^2)^{\frac{1}{2}}} \quad (6.05)$$

and hence

$$r(t) = (r_0^2 + v_{\infty}^2 t^2)^{\frac{1}{2}}$$

A transmission profile measured by X-ray backlighting a homogeneous sphere would show 50% transmission at a radius $r_{\frac{1}{2}}$ given by

$$r_{\frac{1}{2}}(t) = \left\{ r^2(t) - \frac{A^2}{4\rho^2(t)} \right\}^{\frac{1}{2}}$$

where A is the density-length integration needed to give 50% transmission of the appropriate source spectrum. Substituting from equations (6.04) and (6.05) gives

$$r_{\frac{1}{2}}(t) = r_0(1 + \alpha t^2)^{\frac{1}{2}} \left(1 - \frac{\{A(1 + \alpha t^2)\}^2}{2\rho_0 r_0} \right)^{\frac{1}{2}} \text{ where } \alpha = C T_0 / r_0^2.$$

The $r_{\frac{1}{2}}$ trajectories due to the expansion of similar cores but of different initial temperatures can now be computed and compared. In particular, the ratio of experimentally measured $r_{\frac{1}{2}}$ values to simulated (Medusa) $r_{\frac{1}{2}}$ values can be estimated. This is to test the hypothesis that experiment and simulation agree well up to the peak compression phase as

far as core dimensions are concerned but differ afterwards due to hotter cores in the experimental case. Figure 6.16 shows that ratio for parameters typical of shot number 3/27/82 and indicates that the hypothesis is not unreasonable. Parameters used in the calculation were $r_0 = 30 \mu\text{m}$, $\rho_0 = 8 \text{ g cm}^{-2}$, $(C T_0)^{\frac{1}{2}} = 10^6 \text{ cm sec}^{-1}$ for $T_{0s} \approx 20 \text{ eV}$ corresponding to the simulation and the experimental core temperature, (T_{0e}) , was taken in the range 20-300 eV indicated by Abel inversion fits. The data points on the figure from radiographs taken after the implosion indicate that initial core temperatures in excess of the simulation are possible.

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6.3 Stability studies on Ablatively Accelerated Rippled Foils

6.3.1 Introduction

The Rayleigh-Taylor (R-T) instability may represent an obstacle to inertial confinement fusion by disrupting the spherical symmetry of high aspect ratio imploding shells. It is therefore important to obtain quantitative evidence of the growth rate of the instability.

The experiment reported last year (6.01) presented the first clear evidence for the growth of the R-T instability in a laser-accelerated plane target. Subsequently a similar experiment has been carried out and the results indicate reproducibility of the measured growth rate.

6.3.2 Observation of R-T growth

The target was an Aluminium disc of diameter 500 μm and thickness 1.7 μm with a corrugation of wavelength 20 μm and amplitude 0.5 μm . This was symmetrically irradiated with three of the six orthogonal 0.53 μm beams of the 6 beam compression facility. Each of the beams was focused with an f/1 lens at an angle of incidence of 56°, allowing the target normal direction to be used for diagnosis by streaked X-radiography.

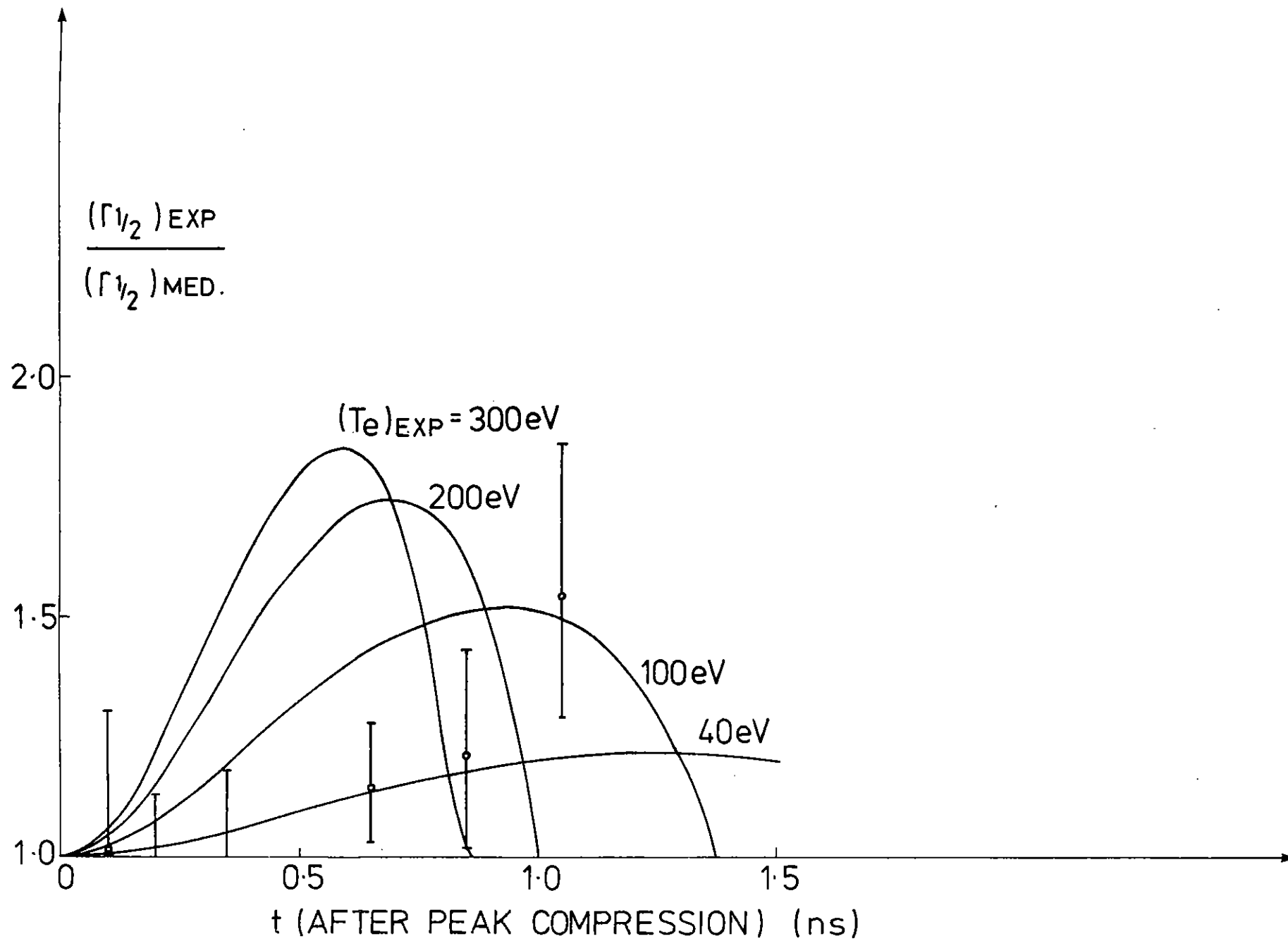


Fig. 6.16 Ratio of experimental and computed core radii for times after peak compression compared to values obtained from a self similar expansion model for various initial temperatures.

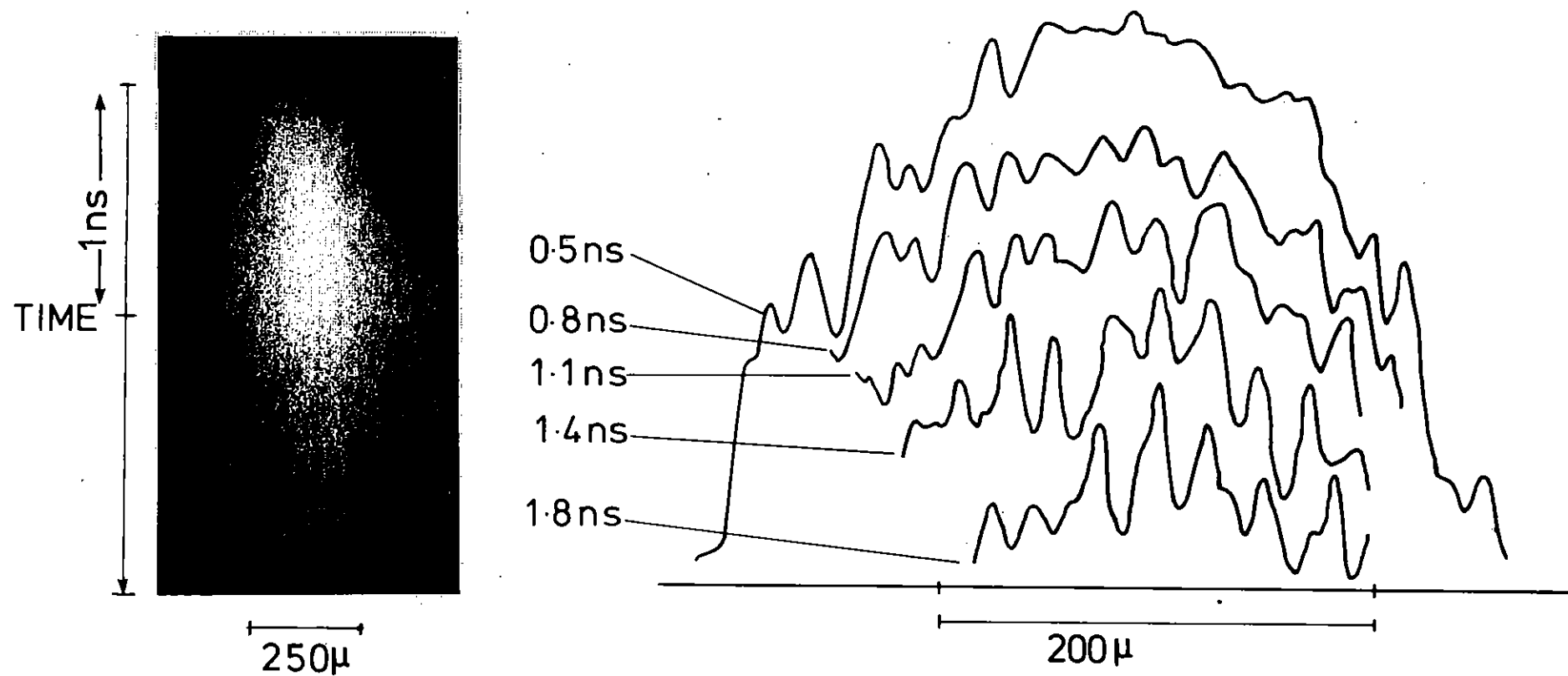


Fig. 6.17 a) X-ray streak photograph of an irradiated microdisc target with 20μm corrugations.
 b) Microdensitometer tracings of (a) at times 500, 800, 1100, 1400 and 1800 psecs, showing increased modulation with time (the density differences between different times are not to scale).

The previous experiment used the X-ray backlighting technique to measure the transmission, and hence mass per unit area modulations in the accelerated target. However, in this case a backlight was not used and the growth of the modulations was observed using the X-ray self-emission of the target. Computer simulation using the 2-D hydrocode "POLLUX" indicates that the source brightness of the self-emission does not vary by more than $\pm 5\%$ throughout the laser pulse, i.e. the front surface of the target acts as its own backlight.

The beams were focused 500 μm beyond the target and the laser energy was 48 J on target with a pulse length of 400 ps. An X-ray streak photograph of the target is shown in Figure 6.17 with the corrugation parallel to the direction of streaking. Initially there is no variation of intensity across the target, but as the instability grows mass is redistributed and modulations in the intensity with a wavelength of 20 μm can be seen.

The amplitude of the modulation of the film density, which is proportional to the areal density modulation is shown in Figure 6.18 as a function of time, the errors representing the spread in the amplitude of the five modulations evident at late time in Figure 6.17. The average growth rate is $2 \pm 0.7 \times 10^9 \text{ s}^{-1}$, with clear evidence of saturation after the end of the laser pulse.

The acceleration of the target was calculated using the measurements of the previous experiment combined with independently measured ablation pressure irradiance scaling relationships (6.02) and allowing for differences in target mass. The corresponding value of the classical growth rate, $(Ka)^{\frac{1}{2}}$, is shown on Figure 6.18. The observed growth rate before saturation has a value of $\approx 0.3 (Ka)^{\frac{1}{2}}$, in agreement with the previous experiment.

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6.4 Analysis of $\lambda = 0.53 \mu\text{m}$ Driven Hot Ablative Implosion Experiments

6.4.1 Introduction.

Laser driven implosion experiments to date have either been in the exploding pusher mode or have been ablatively driven resulting in cooler cores because of the low ablation pressures used. Targets irradiated by $\lambda = 1.05 \mu\text{m}$ laser light have low ablation pressures if the criterion to avoid fast electron preheating, $I\lambda^2 < 10^{14} \text{ W cm}^{-2} \mu\text{m}$, is satisfied. At shorter wavelengths higher ablation pressures, namely $P = 5 (\lambda/1\mu\text{m})^{-2}$ Mbar can be generated at the threshold for fast electron preheating.

Here we describe experiments in which $\lambda = 0.53 \mu\text{m}$ laser light has been used to ablatively compress gas filled glass shells, producing cores hot enough to diagnose by X-ray and thermonuclear product emission. Core temperatures of 200-400 eV are measured. The compressed glass shells had densities of 3 g.cm^{-3} and areal densities of $4 \times 10^{-3} \text{ g.cm}^{-2}$. However, measured compressions are substantially less than predicted by a one dimensional simulation without preheat. The discrepancy is attributed to radiative preheat during the implosion, non uniform irradiation and possibly hydrodynamic instability. Preliminary results from these experiments were reported briefly in last years report (6.03).

6.4.2 Experimental Results

The Laser Division six beam laser system generating 160 J in $\approx 5 \text{ nsec}$ at $\lambda = 0.53 \mu\text{m}$ was used to irradiate glass microballons of 100 to 180 μm diameter and 1 to 2 μm wall thickness filled with Ne (10 bar) or DT (18 bar) + Ne (2 bar), at irradiances from $1 \times 10^{14} \text{ W cm}^{-2}$ to $7 \times 10^{14} \text{ W cm}^{-2}$. The beams were focused 3.6 radii beyond the target centre by f/1 aspheric lenses. Equivalent plane images showed 2:1 irradiance variations on a scale length much less than the beam radius. Six beam ray tracing calculations (6.04) showed 1.2:1 variation in absorbed irradiance with

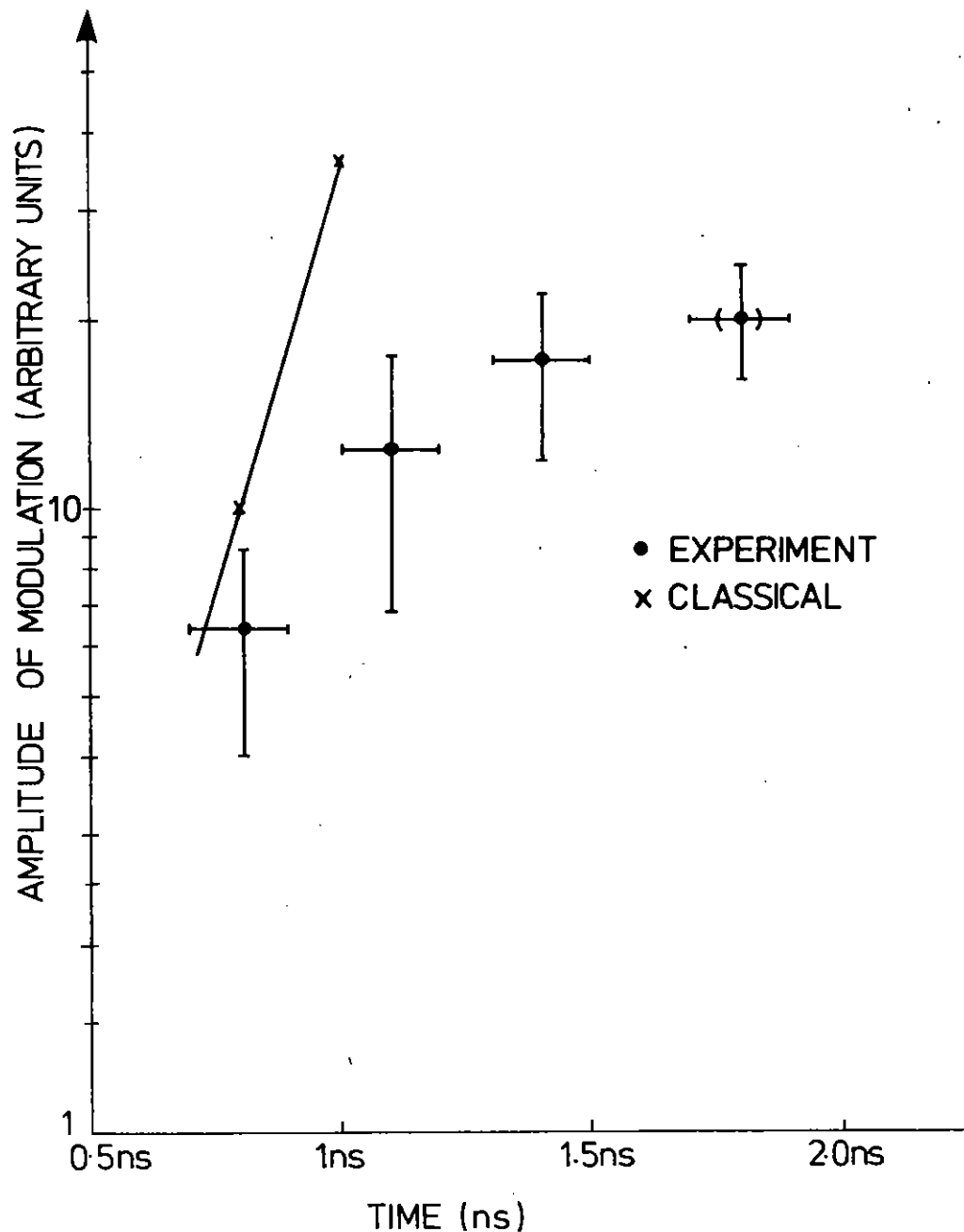


Fig. 6.18 Graph of the amplitude of modulation of film density against time for the latter four measurements in Fig. 6.17. Also shown is the classical growth rate. The absolute amplitude is chosen arbitrarily.

extrema at the centres and triple overlap points.

Fig. 6.19 shows a typical pair of filtered X-ray pinhole camera images in (a) soft ($h\nu \sim 1.5$ keV) and (b) harder ($h\nu \sim 2.4$ keV) X-rays. Emission from the compressed core is seen only in (a) and the observed lower bound on brightness ratio of (a) to (b) implies a temperature of 200 ± 50 eV. By contrast the outer ring of emission from the ablation plasma is of similar brightness in the two images implying a source temperature of 500 eV. The outer ring of emission is a time smeared image of emission from the front and back hot surfaces of the imploding shell. It has a sharp fall in brightness at its inner boundary due to absorption of emission from the back surface by stagnated cool glass shell material in the core. This self radiograph effect gives a measure of the minimum volume occupied by that fraction of the glass shell remaining after the typically 30% mass loss by ablation and thus a density ρ and an areal density ρr (can be calculated. $(3 \pm 1 \text{ g.cm}^{-3}$ and $3 \pm 1 \times 10^{-3} \text{ gcm}^{-2}$ respectively in Figure 6.19). Such low core temperatures and absorption features are not seen in exploding pusher images.

Space resolved X-ray spectra from the core and the outer ring (shell) recorded via a $10 \mu\text{m}$ slit are shown in Figure 6.20. The strong emission from the core at photon energies below the shifted K absorption edge of Si is due to NeIX and NeX recombination continuum. Its slope implies a temperature of 170 ± 20 eV corroborating the estimate from the pinhole camera image. Absorption features ABC in this continuum are due to the surrounding cool glass pusher material, being mainly due to Li and Be like Silicon. Figure 6.20 also shows the Si K absorption edge of this material shifted to $5.65 \text{ \AA}$, consistent with Li to C like ions of Si. Strong SiXIII and SiXIV resonance line emission is mainly from the outer hot shell and masks any underlying absorption features.

Time resolved X-ray spectra were recorded with a novel high brightness focusing spectrometer in which a Von Hamos concave cylindrical pentaerythritol (PET) crystal diffracted a spectrum onto the CsI photocathode of an X-ray streak camera. Figure 6.21 shows a streaked

$\phi = 150 \mu\text{m}$

$\text{D}_2 + \text{T}_2$ 20bar

+

$\Delta r = 1.9 \mu\text{m}$

NEON 1bar

$I = 2.5 \times 10^{14} \text{ Wcm}^{-2}$

3×10^4 NEUTRONS

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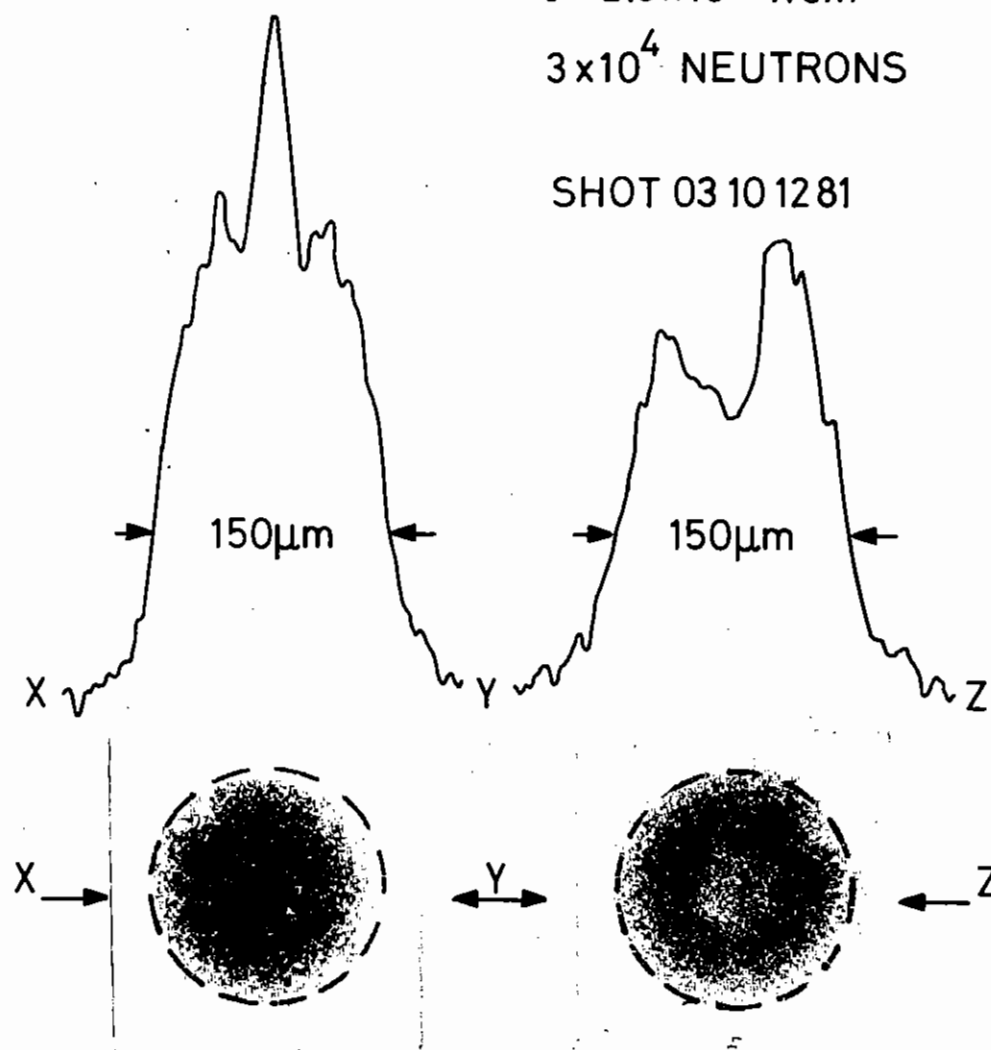


Fig. 6.19 A pair of pinhole images showing soft $h\nu = 1.5 \text{ keV}$ ($10 \mu\text{m}$ pinhole, $25 \mu\text{m}$ Al + $15 \mu\text{m}$ Be filtering) and hard $h\nu = 2.4 \text{ keV}$ ($10 \mu\text{m}$ pinhole, $125 \mu\text{m}$ Be + $10 \mu\text{m}$ Al filtering) images of a typical implosion.

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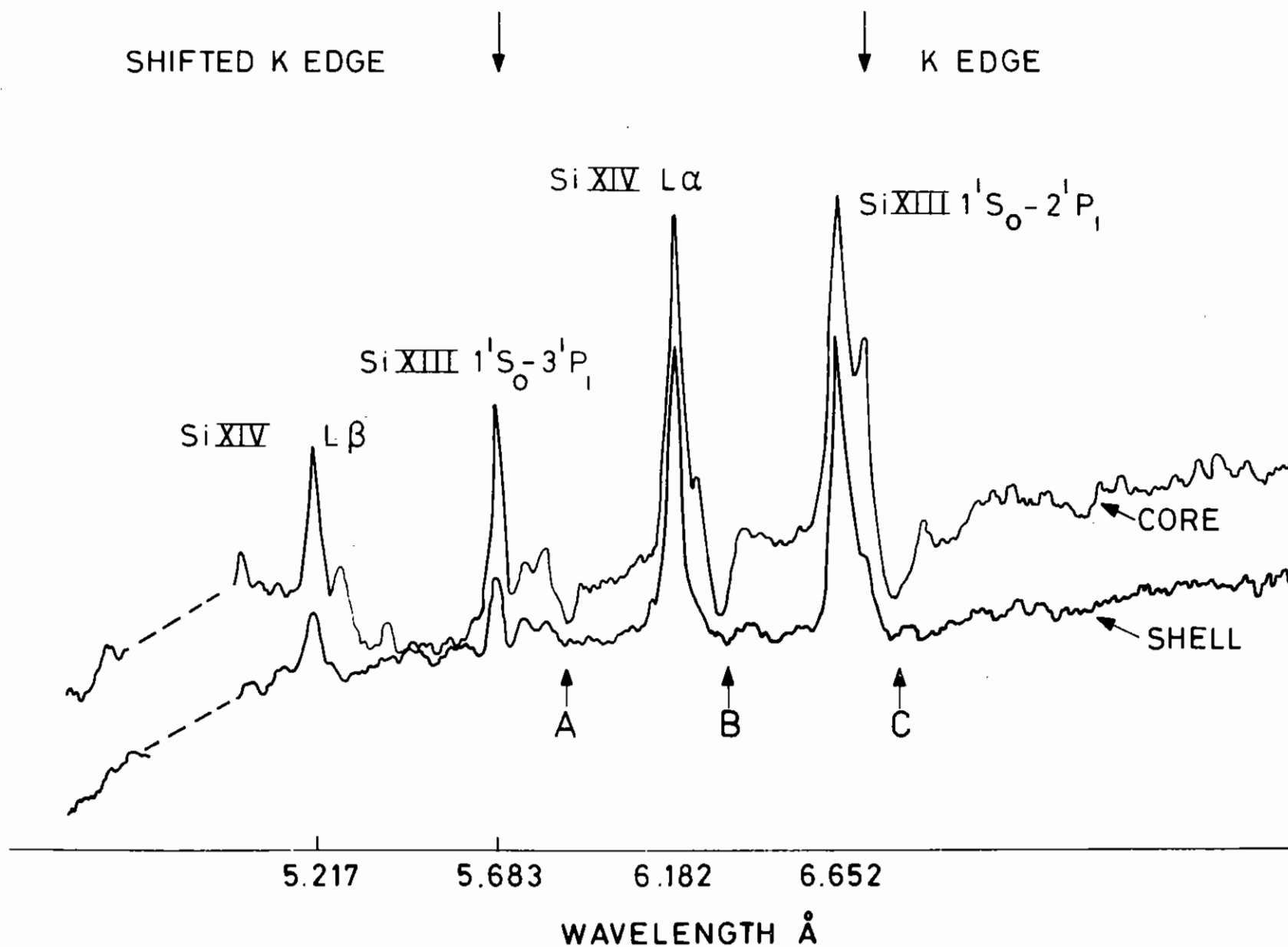


Fig. 6.20 Microdensitometer tracings of a space resolved spectrum (taken with an ammonium dihydrogen phosphate crystal) of a typical implosion.

spectrum in the region of the Si XIV 1s - 2p resonance line. The pulse of satellite line emission (densitometer trace b) follows closely the 500 ps laser power pulse on the glass shell but the adjacent Ne continuum (densitometer trace a) is emitted in a 200 ps pulse whose peak shows that maximum compression is approximately coincident with the peak of the laser pulse.

The ablation pressure P (10 ± 2 Mbar for incident irradiance $1.6 \times 10^{14} \text{ Wcm}^{-2}$ in Figure 6.21) was measured from the implosion time by integrating the equation of motion of the shell. The result is consistent with other measurements of ablation pressure on spherical targets at $\lambda = 0.53 \text{ }\mu\text{m}$ and implies ablation pressures ranging from 7 to 36 Mbar for the 10^{14} to $7 \times 10^{14} \text{ Wcm}^{-2}$ range of irradiance of the present work.

Neutron yields up to 10^5 were recorded from DT filled targets. Using the pinhole camera measurements of the compression the ion temperatures were estimated to be in the range 200 to 350 eV. Ion temperatures were roughly proportional to the estimated specific kinetic energy $pv^2 \sim Pr/\Delta r$ (where Δr is the thickness of the shell).

Alpha particle yields and energy spectra were measured by analysis of etched tracks in CR39 Polymer. Fig 6.22 shows a typical result from the implosion of a thin-walled microballoon ($\Delta r = 1.6 \text{ }\mu\text{m}$) where the total observed yield is 10% of the neutron yield. No alphas could be detected when the microballoon walls were thicker than $1.6 \text{ }\mu\text{m}$. Making due allowance for the tracks falling below the detection threshold, we may estimate that between 0.3 and 0.5 of the α particles produced in the shot illustrated in Fig 6.22 escaped the pusher whose mean pr is thus equivalent to $2.7 \times 10^{-3} \text{ g.cm}^{-2}$ of cold glass or $2.0 \times 10^{-3} \text{ g.cm}^{-2}$ of glass at a temperature of 200 eV. This is consistent with the data from radiography discussed above. The straggling in the observed residual range of a 3.5 MeV α particle is much greater than the resolution of the detector. A 10 to 15% spread in pr can be attributed to simple geometrical effects due to an extended source inside a spherical shell: the additional straggling is presumably due to either time or space-dependent variations in pr .

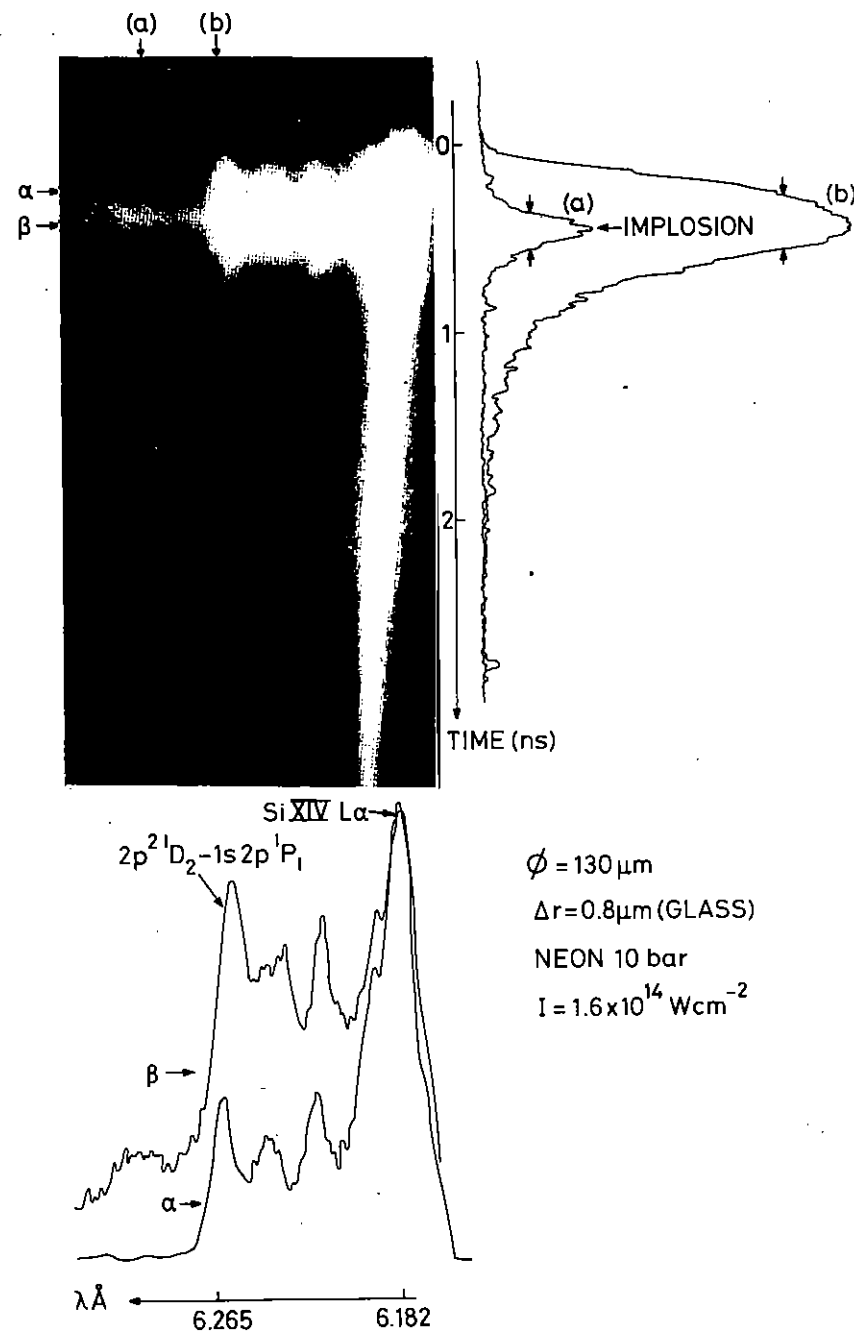


Fig. 6.21 Streaked X-ray spectrum of Si XIV 1s-2p, its dielectronic satellites, and the nearby continuum. The time resolution of the system was 50 psec. An X-ray attenuation of 100 was used to record this spectrum.

One dimensional simulations with negligible hot electron preheating predict values of ρ and ρr in the pusher of 20 g cm^{-3} and $2 \times 10^{-2} \text{ g cm}^{-2}$ respectively with compressed gas temperatures of 2.5 keV. These predicted values are typically an order of magnitude higher than the experimentally observed results.

The simulations did not include preheating by soft X-rays or non spherical deformation of the shell by either Rayleigh Taylor instability or spatial variations in ablation pressure due to non uniform irradiance. These processes all spread the mass of the imploding shell over a wider range of radius, either homogeneously in the case of preheat induced decompression, or inhomogeneously in the case of local deformation, and cause reduced density and temperature in the implosion. Estimates below suggest all three processes are significant in the present experiment although it is more difficult to be sure which is the dominant cause.

The soft X-ray emission from the glass shell was measured in the region $0.9 < h\nu < 1.5 \text{ keV}$ where it agreed with numerical simulation that predicted a spectral flux F_E that peaked in the region of $h\nu \approx 0.5 \text{ keV}$, and an X-ray conversion efficiency of about 5% of the incident energy. Using these values (which are consistent with other measurements of conversion efficiency (Ref. 6.05)), the specific internal energy ϵ of the inside of the shell due to X-ray absorption was estimated to be between 3×10^5 and $3 \times 10^6 \text{ J g}^{-1}$. The uncertainty is due to the effect of transmission near to the oxygen K edge, making the estimate sensitive to the assumed shape of F_E in the 0.5 keV region. The higher estimate of ϵ gives pressures of 50 Mbar at solid density, implying significant in flight decompression of the shell as the ablation pressure is only 10 Mbar.

Numerical simulations with homogeneous preheating at a constant rate predicted densities of the observed magnitude $\rho \approx 2 \text{ g cm}^{-3}$ when preheating reached $3 \times 10^6 \text{ J cm}^{-1}$ but indicated temperatures several times higher than the observed 200 eV. This suggests that other processes are also important, a conclusion reinforced by the observation that with a 1 to 2

μm polymer ablator layer giving about three times less preheat there was no significant difference in the implosions.

Rayleigh Taylor instability growth for wavenumber $k=2\pi/\Delta r$, of 1000 is estimated for these targets assuming a growth rate of about 0.3 of the classical value. This would spread out the mass of the shell and significantly reduce its effective density. It would also mix cool shell material with the hot gas reducing the temperature. One dimensional simulations which include the effects of preheat, indicate that the decompression of the target due to radiation can reduce the density gradients and hence the conditions required for Rayleigh Taylor growth. Two dimensional simulation of the effect of decompression on Rayleigh Taylor growth remains to be carried out.

Spatially non uniform ablation pressure due to the irradiance variations discussed above would also produce a lower density shell and cool core. Thermal smoothing is quite effective for nanosecond pulses at the typically $4 \times 10^{14} \text{ W cm}^{-2}$ irradiance of this experiment but is much reduced by the finite build up time of the separation between ablation front and absorption region on the 250 psec time scale of these implosions. Variations of up to 1.2:1 in distance travelled by the imploding shell are therefore expected with an associated reduction of the mass at any radius by an order of magnitude.

It seems that the only way to avoid these problems is to use shells of lower aspect ratio $r/\Delta r$. For a given radius the increased Δr dramatically reduces X-ray preheat. The increased mass lengthens the implosion time and improves thermal smoothing. Also Rayleigh Taylor growth, which scales as $\exp(r/\Delta r)^{1/2}$ in the low preheat limit, is reduced.

Estimates suggest that with $r \approx 100 \mu\text{m}$ and $r/\Delta r \approx 15$, implosions approaching the ideal ablative mode may be obtained with temperature $kT \approx 1 \text{ keV}$ and density up to 20 g cm^{-3} suitable for emission based diagnostics. There is a minimum radius for shells of a given aspect ratio imposed by the need to limit soft X-ray preheating and thus a minimum laser energy of the order of 0.5 kJ.

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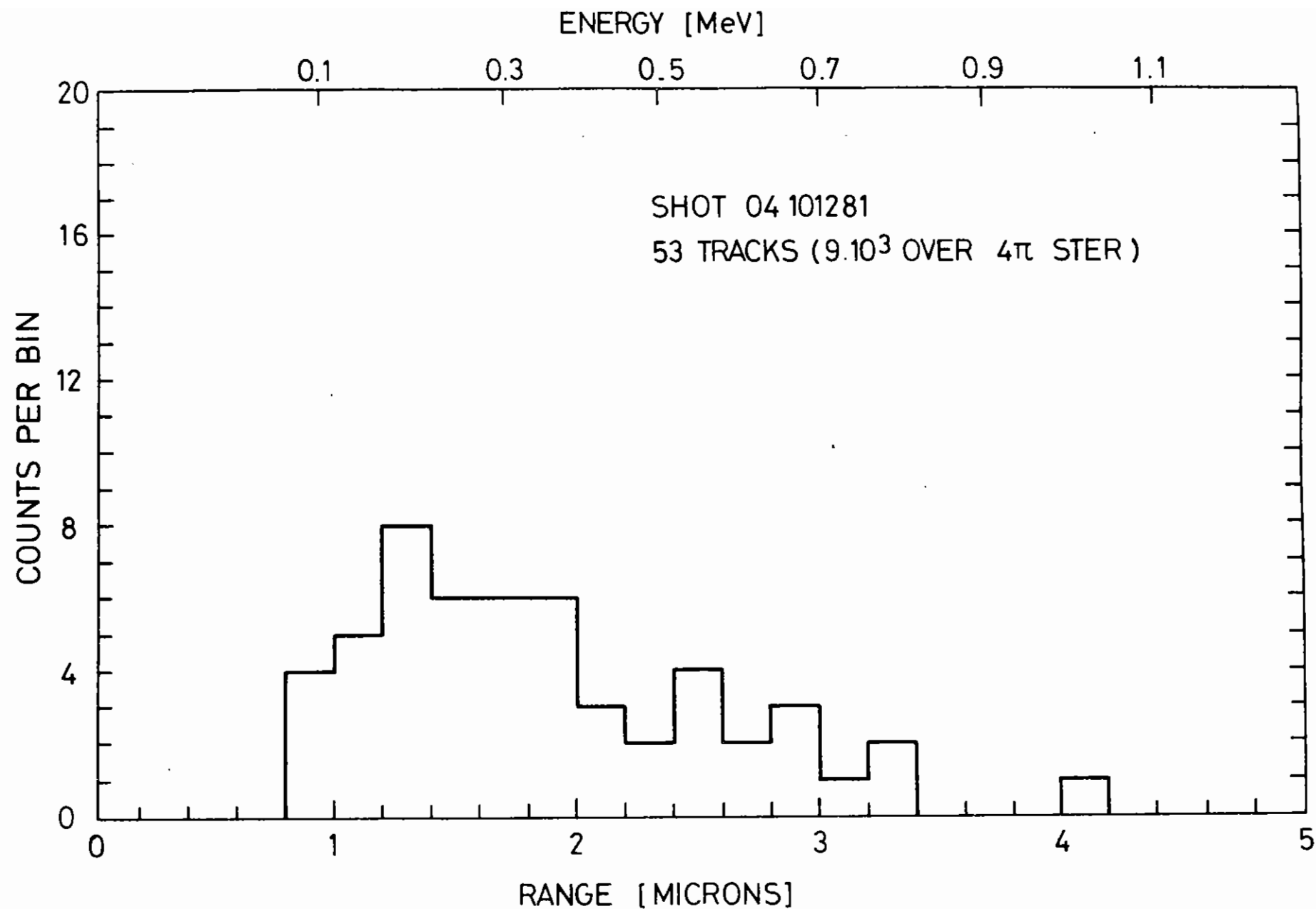


Fig. 6.22 Histogram of alpha particle track lengths detected in CR39 polymer. The corresponding incident alpha particle energy is shown. The neutron yields were a factor of 10 higher than the measured number of alpha particles escaping from the target.

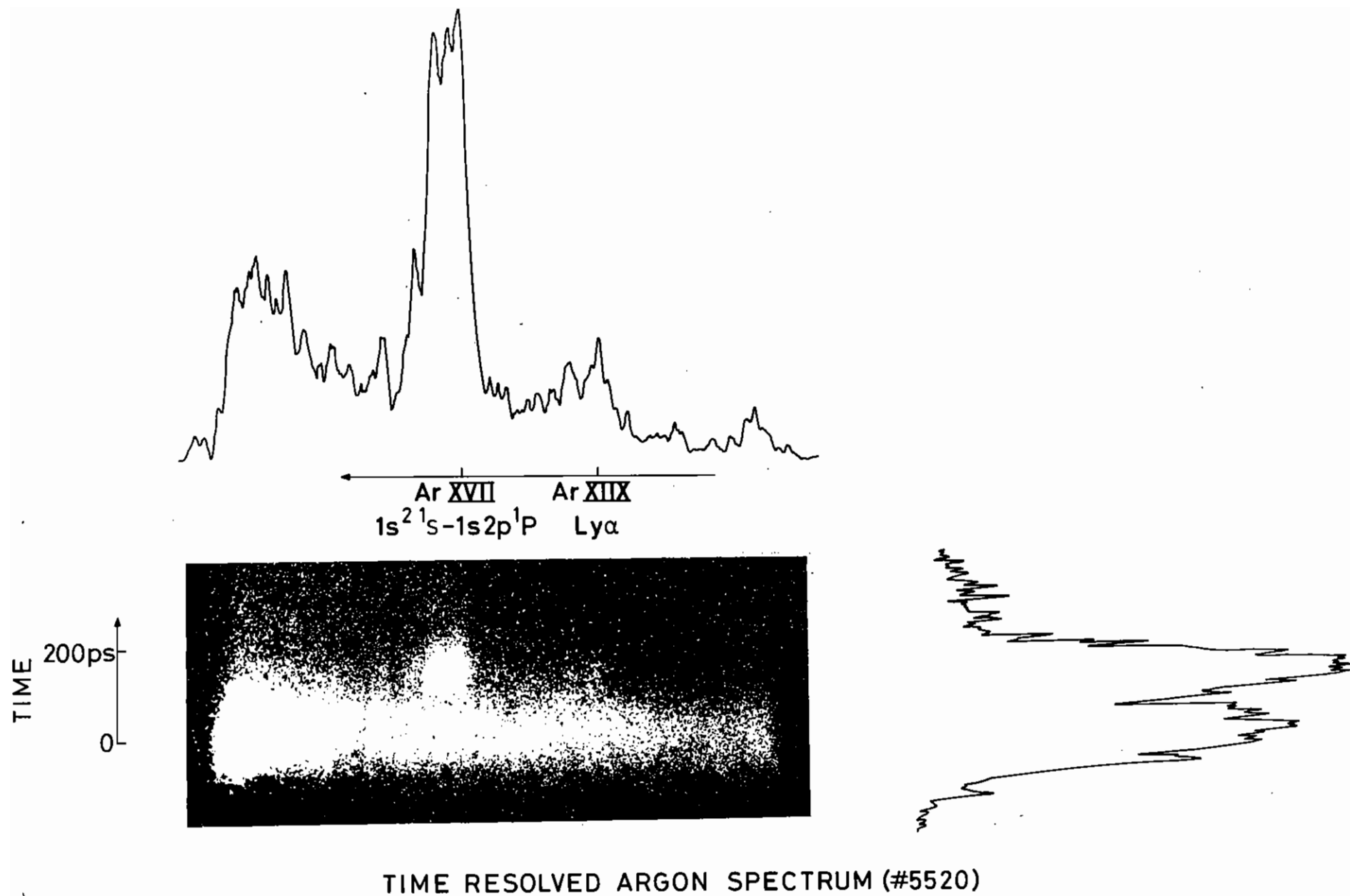


Fig. 6.23 Time resolved Argon spectrum for shot number 5520.

6.5.1 Introduction

In order to continue the $\lambda = 0.53$ implosion experiments described in Section 6.4, a collaborative experiment was carried out at KMS Fusion Inc using X-ray diagnostic equipment provided by RAL. The experiment was performed using the spherically symmetric (TBIS) illumination system at KMS in order to achieve greater target irradiation uniformity than that possible with 6 beams at RAL. Unfortunately due to limitations caused by mirror damage in the system it was not possible to carry out long pulse (~ 1 ns) experiments. Therefore the experiment was performed using pulses of 100 (or 200) ps duration at incident irradiances of up to $10^{15} \text{ W cm}^{-2}$. Since target shell thicknesses were of order $1 \mu\text{m}$ these implosions were therefore dominated by fast electrons ($I\lambda^2 \approx 2-3 \times 10^{14} \text{ W cm}^{-2} \mu\text{m}^2$) and hence were 'quasi - exploding pusher' implosions.

6.5.2 Experiment and diagnostics

Glass microshells of 70 - 150 μm diameter were irradiated with up to $10^{15} \text{ W cm}^{-2}$ in 100 - 200 picoseconds. The Triple Bounce Illumination System (TBIS) uses two hemispherical mirrors to focus light onto the target after it has experienced three reflections. The illumination is then incident on the target from almost 4π steradians. A narrow equatorial gap between the mirrors allows diagnostic access to the target, the absence of illumination in this equatorial region is prevented by defocusing the mirrors slightly so that rays overlap on the equator.

Three different classes of target were imploded and found to emit seed gas spectra.

- (i) Small (70 μm diameter), thin-walled (0.9 μm) glass microballoons filled with 0.4 atmospheres of Argon.
- (ii) Larger ($\sim 110 \mu\text{m}$ diameter) glass microballoons with wall thicknesses between 0.9 μm and 1.5 μm filled with 2 atmospheres of Argon.
- (iii) Microballoons as in (ii) but with an extra fill of Deuterium gas (D_2) of the order of 100 atmospheres.

X-ray diagnostics employed in this experiment included time-resolved and time-integrated spectroscopy. Time-resolved spectra were recorded using a flat RbAP crystal (Rubidium acid phthalate, $2d = 26.12 \text{ \AA}$) to diffract X-rays onto the photo-cathode of an X-ray streak camera (on loan from Rutherford Appleton Laboratory). Figure 6.23 is a typical streaked Argon spectrum. The time resolution is 40 picoseconds, sufficient to show the delay between the bright continuum, corresponding to the laser ablating the shell, and the appearance of Argon lines. The Argon $1s^2 \text{ } ^1\text{S}_0 - 1s2p \text{ } ^1\text{P}_1$ ($\text{He}\alpha$) and $\text{Ly}\alpha$ are seen on this streak.

Time-integrated spectra were observed using a minispectrometer containing a PET crystal (Pentaerythritol, $2d = 8.74 \text{ \AA}$). The spectrometer was positioned with its entrance slit 6 cm from the target instead of its usual position 5 mm from the target. The large separation between spectrometer and target (necessary due to the TBIS geometry) meant that space-resolved spectroscopy was not possible. However the larger separation meant higher dispersion spectra, at the expense of reduced spectral range. A typical Argon spectrum is shown in Fig. 6.24.

The electron number density n_e in the compressed core was calculated from the Stark-broadened line profiles of Argon $\text{He}\beta$ and $\text{He}\gamma$. The profiles were taken from time-integrated spectra such as that of Fig. 6.24 and digitised. The line profiles were then fitted in n_e , n_l & parameter space to the optically-thin Helium-like profiles of R.W. Lee with L.T.E. opacity broadening (6.06). Here n_l is the line density of Helium-like Argon ions (Ar XVIII) in the ground state contributing to the opacity broadening of the line. Electron densities were found to be in the range

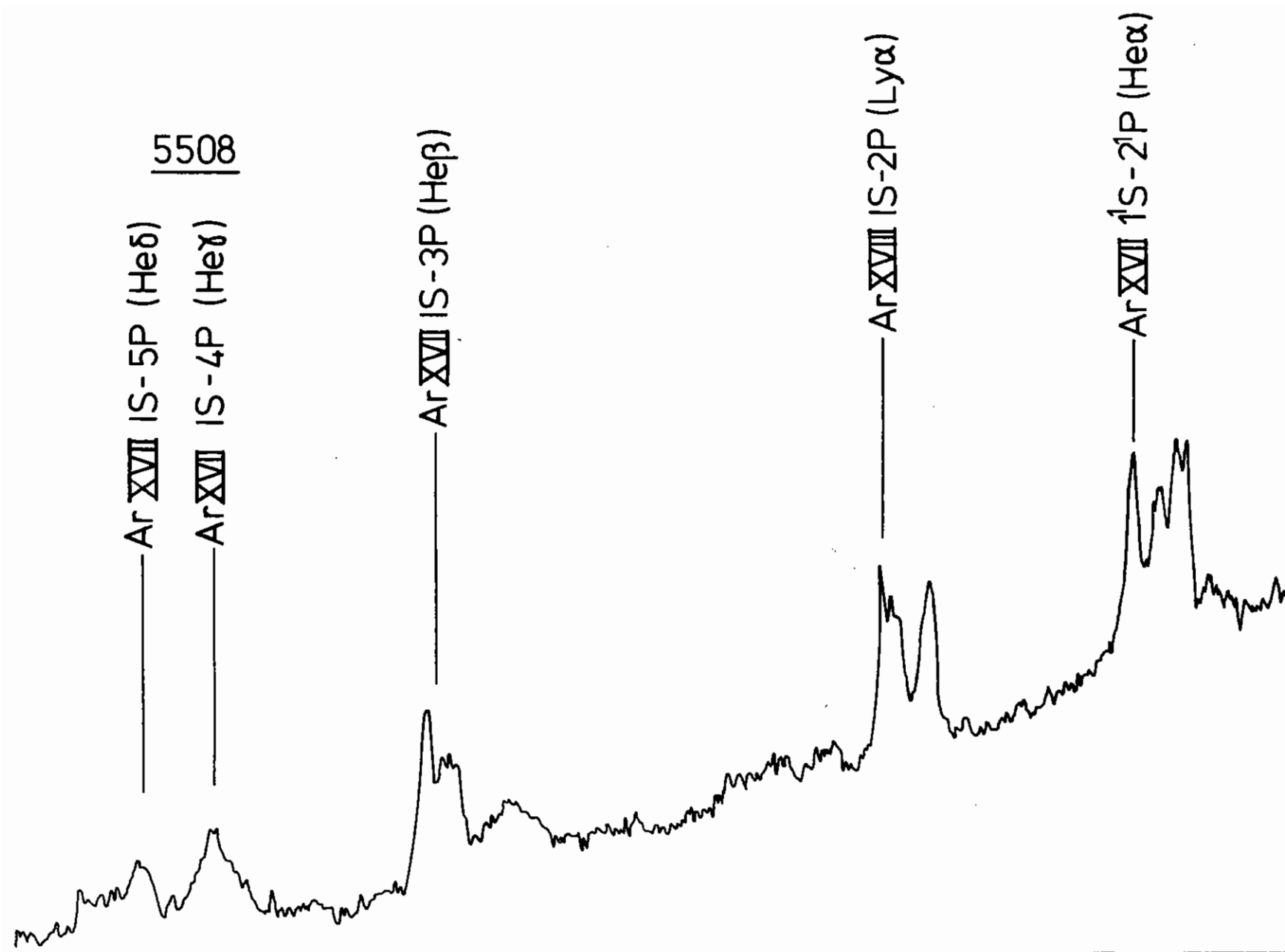


Fig. 6.24 A typical time integrated Argon spectrum from a KMS green implosion. Density diagnosis was possible through line shape analysis. Temperature determination was achieved by comparison of He - like and H - like line intensities.

10^{23} cm^{-3} to $3 \times 10^{23} \text{ cm}^{-3}$ equivalent to mass densities between 0.3 and 1 g. cm^{-3} .

The ratio of Helium-like to Hydrogen-like Argon lines served as a diagnostic of electron temperature in the compressed core. The relative intensities of the different lines were calculated using the population calculations of R.W.Lee. The collisional and radiative rate equations are solved for a homogeneous plasma and the optically thin line ratios calculated. The collisional excitation rates were updated to agree with the calculations of Pradhan et al. and Lee (6.07, 6.08). The results of the calculations are summarised in Figure 6.25 where various line ratios are plotted as a function of electron temperature for two different electron number densities. The ratio used for temperature diagnosis in this experiment was that of the Ly α to He γ . The ratio is seen in Fig. 6.25 to be rapidly varying with temperature in the range 500 - 900 eV and insensitive to electron density changes.

6.5.3 Results

The electron densities and electron temperatures observed are summarised in Table 6.02. The specific absorbed energy quoted is obtained by dividing the absorbed energy, as measured by an array of four plasma calorimeters, by the initial mass of the target. The error quoted for the electron temperature is from the uncertainty in measuring the line ratios. The final column in Table 6.02 gives a qualitative measure of the brightness of the time-integrated spectra obtained. The brightness quoted is proportional to the intensity of the group of lines near the Argon He α line that includes the $1s^2 \ ^1S_0 - 1s2p \ ^1P_1$ resonance line and $1s^2 \ ^1S_0 - 1s2p \ ^3P$ intercombination line.

The brightness of the spectra is a function of the shell thickness. The electron temperature is also a function of the shell thickness but varying the shell thickness also affects the specific absorbed energy. The final temperature in these quasi-exploding pusher implosions should be linearly dependent on the absorbed specific energy. Figure 6.26 shows the electron temperatures plotted as a function of specific absorbed energy for the

shots in Table 6.02. The target type in the key of Fig. 6.26 refers to the description in Section 6.5.2.

There is an obvious correlation between temperature and specific absorbed energy. Comparing shot 5508 with 5509 (in Table 6.02) which have similar electron temperatures and densities, the brightness is seen to be lower by a factor of three in the latter shot. The target diameter and Argon fill density were similar for these two shots but the shell thickness was greater for 5509. A possible explanation for the reduced brightness in the spectrum of shot 5509 is that the higher compressed shell pressure, expected due to the greater initial shell thickness, is attenuating the Argon emission. The transmission of glass (SiO_2) as a function of pusher pressure and temperature is plotted in Fig. 6.27 for an X-ray wavelength of 3.96 Å (Argon He α). The possibly lower pusher temperature associated with the thicker shell of 5509 coupled with a higher expected pressure, indicates that the glass transmission could be significantly different between shot 5508 and 5509. The region of rapid variation of shell transmission with pressure of $2 \times 10^{-3} \text{ g cm}^{-2}$. Hence in the ablative implosions described in Section 6.4 any Argon lines excited in the core would be severely attenuated by the high line density of the pusher ($3 - 6 \times 10^{-3} \text{ g cm}^{-2}$).

Shot 5523, which was of similar diameter to 5508 and 5509 but had a wall thickness of 1.4 μm , produced very weak spectra (see Table 6.02). The core temperature in this implosion was possibly lower but the weakness of the Argon lines observed meant that the electron temperature determination was less accurate.

The electron temperatures observed with D_2 filled targets were lower than for pure Argon filled targets but the specific absorbed energies were also lower (see Figure 6.26).

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TABLE 6.02

	τ_{probe} w.r.t. τ_{imp}	τ_{imp} w.r.t. Peak of the Driver pulse	Core Radius at $\approx 50\%$ Transmission Point		ABEL FITS		Medusa		ABEL FIT	Medusa
Shot no., Date	$\tau_p(\text{exp})$ (+) ns	$\tau_p(\text{Med})$ ns	$r_{\frac{1}{2}}(\text{exp})$ μm	$r_{\frac{1}{2}}(\text{Med at } \tau_p)$ μm	Peak Dens. Abel Inv. $\text{g}\cdot\text{cm}^{-3}$	Av.dens. $\text{g}\cdot\text{cm}^{-3}$	Peak Dens. $\text{g}\cdot\text{cm}^{-3}$	Av.dens. $\text{g}\cdot\text{cm}^{-3}$	(Te)av. eV	(Te)av. eV
05,26/8/82	-0.85	0.85	49.5 $\pm$ 3.5	50.2 $\pm$ 3.6			3.2 $\pm$ 0.1	1.6 $\pm$ 0.1		2.1
08,26/8/82	-0.65	0.65	46.6 $\pm$ 3.0	45.2 $\pm$ 3.2			3.4 $\pm$ 0.2	1.7 $\pm$ 0.1		2.3
10,26/8/82	-0.3	0.5	35.0 $\pm$ 1.5	31.3 $\pm$ 2.3	1.4 $\pm$ 0.3	0.7 $\pm$ 0.2	3.7 $\pm$ 0.2	0.9 $\pm$ 0.2	90 $\pm$ 5	2.6
17,26/8/82	-0.75	0.95	51.7 $\pm$ 1.5	51 $\pm$ 2.5			3.2 $\pm$ 1.03	1.6 $\pm$ 0.1		2.1
21,26/8/82	-0.05	0.45	21 $\pm$ 2.2	18.0 $\pm$ 2.0	4.7 $\pm$ 1.1	2.5 $\pm$ 0.5	24.8 $\pm$ 1.2	3.5 $\pm$ 0.8	160 $\pm$ 10	12.0
01,27/8/82	0.1	0.3	18.1 $\pm$ 1.4	17.9 $\pm$ 2.9	6.5 $\pm$ 1.5	3.9 $\pm$ 0.8	10 $\pm$ 5	3.1 $\pm$ 1.6	185 $\pm$ 20	13.2
03,27/8/82	0.2	0.4	18.8 $\pm$ 1.6	21.6 $\pm$ 3.6	5.6 $\pm$ 1.4	3.6 $\pm$ 0.8	9.5 $\pm$ 4.0	2.5 $\pm$ 1.2	200 $\pm$ 30	6.5
06,27/8/82	0.35	0.35	22.4 $\pm$ 2.2	24 $\pm$ 3.1	3.3 $\pm$ 0.7	1.7 $\pm$ 0.3	4.6 $\pm$ 0.8	1.2 $\pm$ 0.2	135 $\pm$ 10	2.7
08,27/8/82	-0.1	0.8	56.4 $\pm$ 2	25.3 $\pm$ 3.0			23.3 $\pm$ 16.8	1.9 $\pm$ 0.5		5.1
07,2/9/82	-1.05	0.75	72.1 $\pm$ 4.0	60.7 $\pm$ 3.1			3.0 $\pm$ 0.01	1.5 $\pm$ 0.5		1.8
10,2/9/82	-0.85	0.35	43.3 $\pm$ 4.3	51.3 $\pm$ 3.0			3.5 $\pm$ 0.1	1.8 $\pm$ 0.1		2.2
15,2/9/82	0.65	0.65	28.7 $\pm$ 1.8	33.8 $\pm$ 2.2	2.7 $\pm$ 0.5	0.6 $\pm$ 0.1	1.9 $\pm$ 0.4	0.69 $\pm$ 0.1	40 $\pm$ 5	1.5
17,2/8/82	0.85	0.75	41.6 $\pm$ 4.0	34.3 $\pm$ 2.5	1.9 $\pm$ 0.4	0.6 $\pm$ 0.1	1.2 $\pm$ 0.3	0.7 $\pm$ 0.1	80 $\pm$ 10	0.8
19,2/9/82	1.05	0.75	53.75 $\pm$ 2.0	35 $\pm$ 5.0	1.0 $\pm$ 0.1	0.4 $\pm$ 0.1	1.0 $\pm$ 0.1	0.9 $\pm$ 0.36	30 $\pm$ 5	1.4

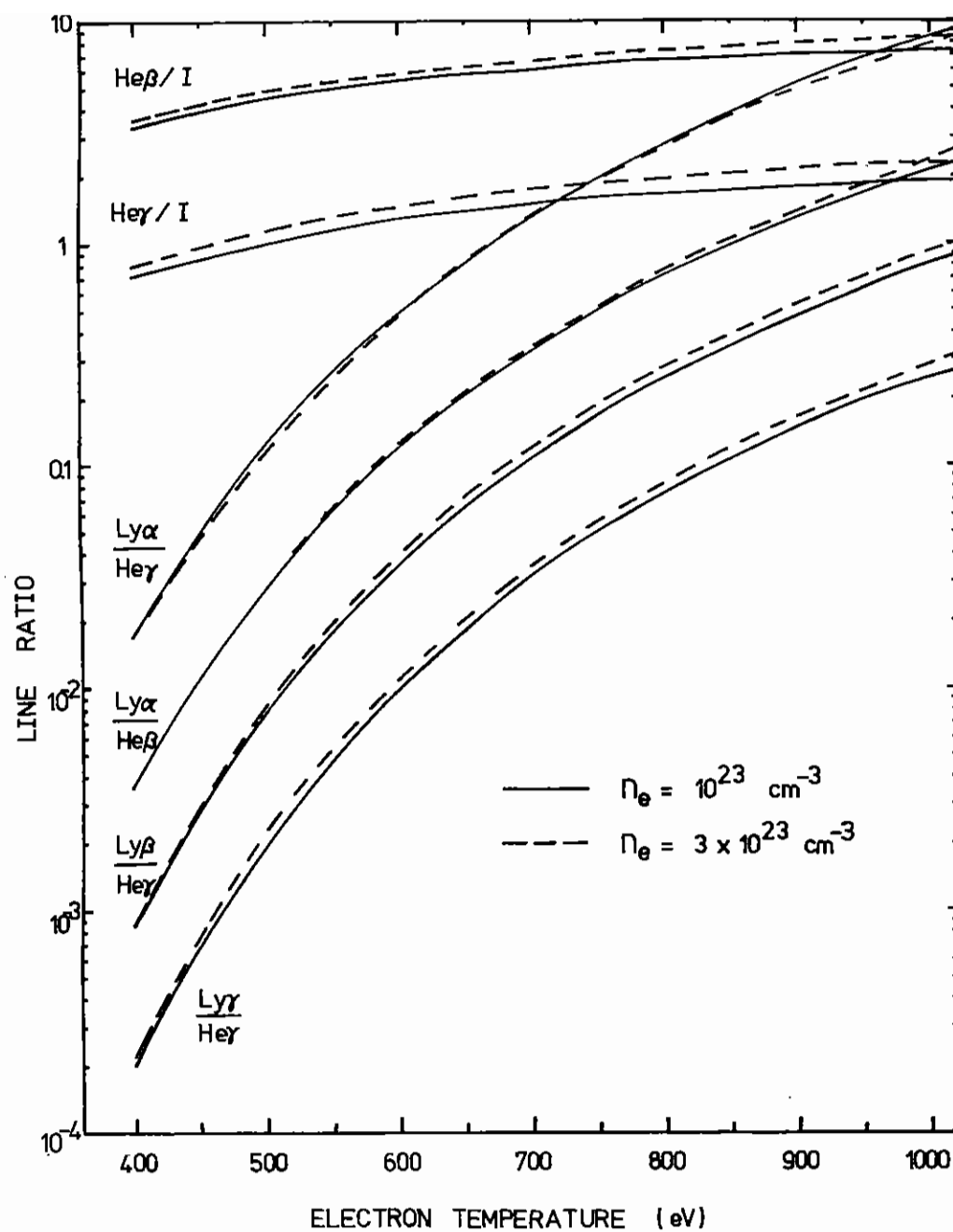


Fig. 6.25 Ratios of the intensities of the main X-ray lines of H-like He-like Argon as a function of electron temperature for two different electron densities. The intercombination line, $1s^2 \ 1S_0 - 1s \ 2p^3P$ is referred to as 'I'.

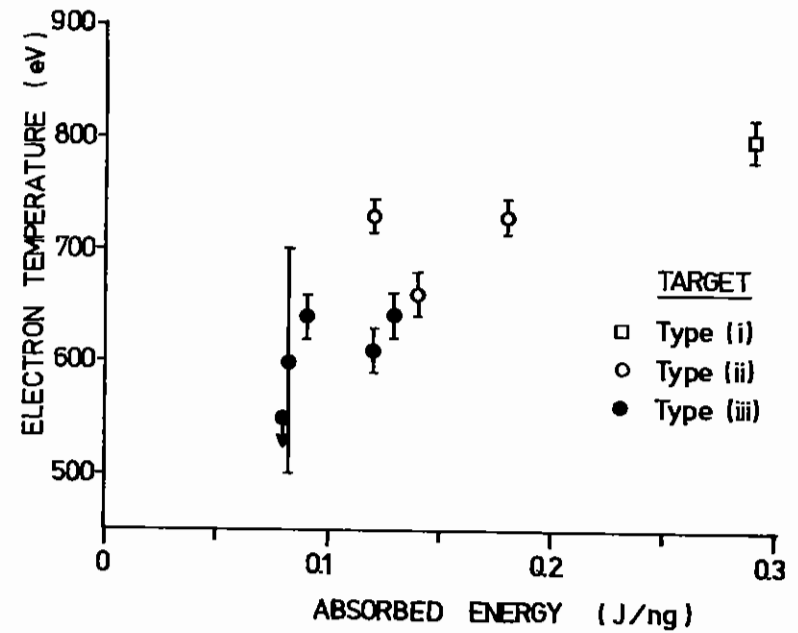


Fig. 6.26 Measured core electron temperatures from Table 6.03 plotted against specific absorbed energy. The target types are described in the text.

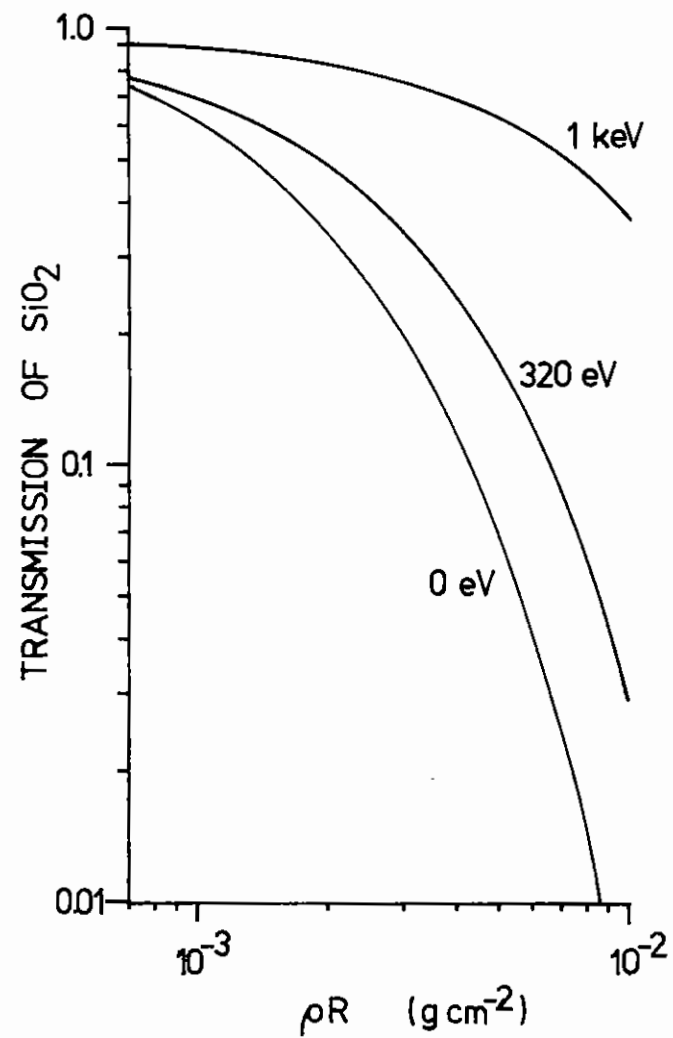


Fig. 6.27 Transmission of glass pusher as a function of line density for three different glass temperatures (at $\lambda = 3.96 \text{ \AA}$).

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INDEX

7.1	INTRODUCTION	page 7.1
7.2	CVI EXPERIMENT	page 7.2
7.2.1	Introduction	
7.2.2	Experiment with Single Green Beam Illumination	
7.2.3	Experiment with Multiple Green Beam Illumination	
7.2.4	Data and Analysis	
7.2.5	Computer Modelling of Fibre Under Green Beam Illumination	
7.3	DIELECTRONIC SATELLITE SPECTRA	page 7.11
7.3.1	Density-dependent Si XII Dielectronic Satellite Spectra	
7.3.2	The Density Dependence of Dielectronic Satellites of Helium-like Transitions	
7.4	K-EDGE FIDUCIALS TO SEE ABSORPTION SHIFT IN RESONANCE LINE POSITIONS	page 7.24
7.5	QUANTITATIVE X-RAY MEASUREMENTS FOR RADIATIVE PREHEAT DETERMINATION	page 7.24
7.6	TIME DEPENDENT STARK BROADENING MEASUREMENTS IN PLANAR LASER TARGET EXPERIMENT	page 7.26
7.6.1	Introduction	
7.6.2	Experimental Design	
7.6.3	Theoretical Modelling	
7.6.4	Results	
7.6.5	Discussion	
7.6.6	Conclusion	

7.7	STUDYING ELECTRON DISTRIBUTIONS USING THE TIME-RESOLVED FREE-BOUND SPECTRA FROM CORONAL PLASMAS	page 7.36

7.7.1	Introduction
7.7.2	Definition of Experiment
7.7.3	Experimental Method
7.7.4	Discussion
7.7.5	Conclusions

7.8	TIME RESOLVED X-RAY LINE DIAGNOSTICS OF LASER-PRODUCED PLASMAS	page 7.43

7.9	EXAFS SPECTROSCOPY	page 7.46

7.9.1	Introduction
7.9.2	Experiment
7.9.3	Time-Resolved EXAFS

REFERENCES	page 7.52

CHAPTER EDITOR: R W Eason

7.1 Introduction

In the past year the effort of the group has been directed into two main areas. XUV laser development and XUV and x-ray spectroscopy as a diagnostic tool. It is expected that this pattern will continue in the future. However, as experimental techniques allow it is anticipated that the efforts of the group will also be directed to investigating the properties of high density plasmas developed by compression.

Specific experiments in the diagnostic field have been concerned both with plasma and neutral states. The analysis and measurements of dielectronic satellite spectra has been investigated as a measure of plasma parameters in hot dense media. The problem of electron density measurements by Stark broadening in such plasmas has been studied experimentally to compare with detailed calculations of such line profiles. In the hot tenuous plasmas of the corona free-bound recombination radiation has been used to measure electron distribution functions.

In the atomic state x-ray absorption spectroscopy has been used to examine behaviour near the ionisation edge, in particular the occurrence of EXAFS, and density induced line shifts. In a similar experiment the role of pre-heat in the target interior from a surface laser plasma was examined.

Work has proceeded on the carbon fibre XUV laser programme. The major steps have been technological involving the use of 0.53 μ m green light and four beam irradiation. Significant gain has been measured, but the system is at present not optimised. Computer modelling is in qualitative agreement with the experiments, and should allow future work to be directed towards higher gain conditions.

7.2 CVI Experiment7.2.1 Introduction

The development of the working XUV laser based on the carbon fibre expansion-recombination scheme is perhaps the most likely approach at present in this field. However before the proof of principle experiments performed at Hull can be realised as a working laser system, important technological modifications must be made to the pumping configuration. Since the Hull experiments were concerned only with a demonstration of a useable gain coefficient, sufficient in principle to construct a laser, the experimental arrangement could be devised in a simple manner not necessarily readily scaled. To this end a simple single beam, one-sided illumination system using fundamental Nd: glass radiation was used (7.01). In a working laser attention must be given to technological constraints of beam focusing, fibre alignment etc., to construct a system scalable to a sufficiently, large gain-length product.

Such design considerations based on currently available lens systems make one sided illumination extremely complex. A more attractive technological solution using multi-beam illumination is a symmetric lens arrangement. In such a system simple microscope alignment as used at Hull becomes unpracticable and techniques such as those already used at the Central Laser Facility must be employed (7.02). With thin fibres such methods are considerably more accurate if used with visible green radiation, rather than infrared. The technological constraints thus indicate the use of 0.53 μ m radiations with its attendant advantages in laser-fibre coupling, higher flux limiting etc.

The experiments carried out over the past year have been performed with these constraints in mind. The role of green light as a target illuminant was studied initially in single beam, one sided illumination, and later in a four beam experiment. In the latter series, gain was measured in four runs, and an indication of gain in several more. However more importantly the experiments correlate well with theory and allow us to identify the problems associated with the present arrangement. As will be discussed

in a subsequent section, these principally concern the matching of the fibre characteristics to the laser. Three solutions may be suggested.

- 1) Use of thinner fibres (3-4 μm diameter)
- 2) Improved heating symmetry
- 3) Modified laser pulse characteristics

Of these the use of thinner fibres is the most straight forward, but may be subject to technological difficulty: the best option has not as yet been fully examined.

7.2.2 Experiment with single green beam illumination

The CVI experiments described in the last Annual Report (7.02) were designed to characterize the energy coupling of a $1.06\mu\text{m}$ ~ 100 ps pulse into a ~ 2mm length of $5\mu\text{m}$ diameter carbon fibre. It was found that the average energy coupling was ~ 5.5% with a peak value ~ 9.5%. The alignment technique using a $1\mu\text{m}$ laser was extremely difficult and time consuming, suggesting that a full scale gain experiment at the Central Laser Facility would have to be done using the second harmonic. An experiment using four green beam illumination was planned but thought to be rather ambitious to attempt in a simple step, so an intermediate, single green beam experiment was carried out first. From this experiment a direct comparison with earlier $1\mu\text{m}$ results could be obtained, and it was hoped that some optimization of the laser characteristics would also be possible. Additionally a comparison of the spectrograph sensitivities would allow an appropriate choice to be made for the four beam experiment.

7.2.2(a) Energy coupling data

Two calorimeter arrays were used throughout the experiments: a six element and a four element. The six element array was kept fixed at a position 50° below the North lens. The four element array was moved to various positions around the fibre axis so that the assumption of cylindrical symmetry, used in the calculation of the coupling, could be

checked. Both calorimeters were ~ 20 cm from the target. After allowing for the different sensitivities and applying appropriate corrections to the signals the symmetry was demonstrated to be very good (Fig. 7.01). In this figure the solid points are actual measurements whilst the open points have been obtained by invoking symmetry along the north-south line. The point which lies directly on the north-south line was obtained by placing a single calorimeter element in the centre of the focusing (north) lens. In all cases the peak signal from the six element array has been used for normalization. The good symmetry is in agreement with space-resolved measurements at Hull (7.01).

The initial experiments used a pulse length ~ 100 ps with a 20% prepulse 200 ps before the main pulse. The average coupling was ~ 5.5% with a peak ~ 9.5% and so is comparable with the measurements using the fundamental. The fact that the second harmonic does not give an appreciably different coupling is not particularly surprising since the energy absorption is strongly affected by the size of the critical density surface of the prepulse plasma, which is smaller for the second harmonic.

Code modelling does suggest, however, that a higher coupling can be achieved with a longer pulse length, with ~ 200 ps being optimum. This is confirmed by the measurements made using a 200 ps pulse which show an average coupling ~ 9% with a peak ~ 17%: a 20% prepulse and 200 ps delay was used again. Similar measurements using ~ 160 ps pulses at the fundamental wavelength have also shown improved coupling over 100 ps pulses in experiments at Hull. Fig. 7.01 showed the symmetry of the expansion and this is once again highlighted in Fig. 7.02, which shows calorimeter measurements for arrays almost diametrically opposite each other. The similarity between the profiles is quite evident and both measurements indicate ~ 14.5% coupling in this case.

Increasing the prepulse/main pulse delay to 500 ps for some shots had no appreciable effect on the coupling confirming the conclusions from the measurements at the fundamental wavelength that the delay is not critical (7.02). Whether the delay affects the gain as little as the coupling is less certain and requires further work.

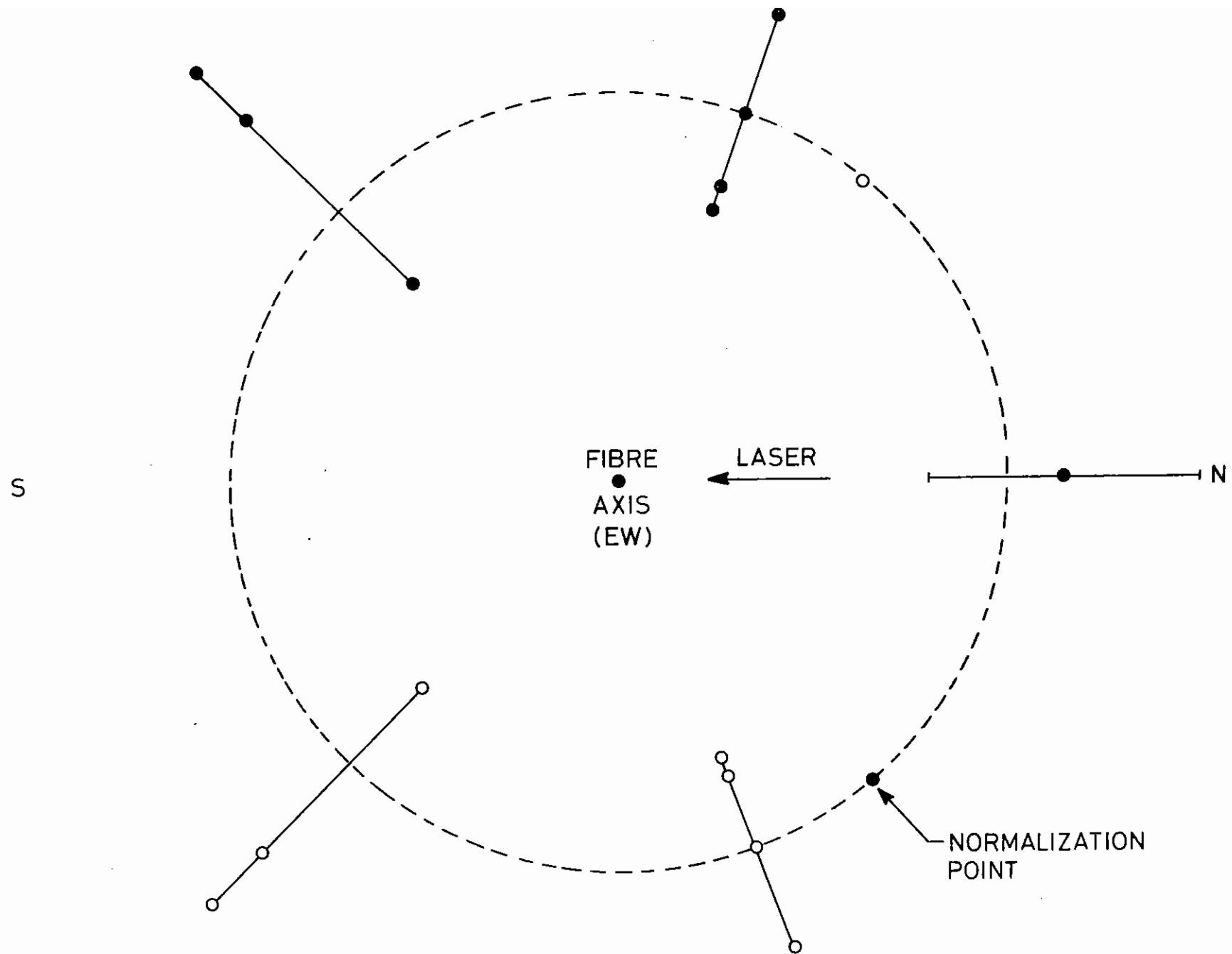


Fig. 7.01 Demonstration of cylindrical symmetry: solid points are actual measurements, and the open points have been obtained by invoking symmetry along the N-S line.

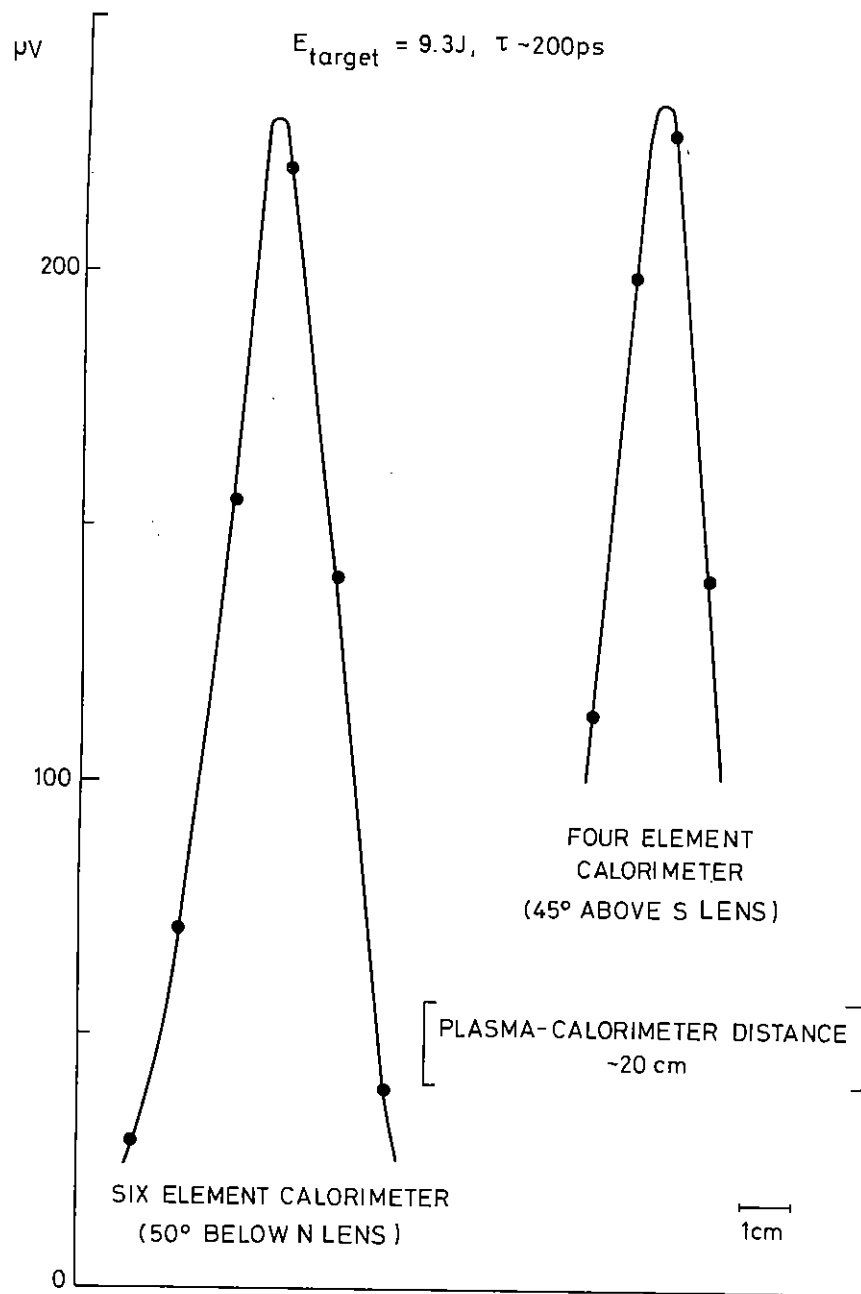


Fig. 7.02 Calorimeter measurements for arrays almost diametrically opposite each other. The similarity between the profiles is evident. Both measurements indicate ~14.5% coupling in this case.

7.2.2(b) Spectrograph comparison

Throughout the experiment the fibre plasma was viewed axially by a compact 1 m grazing incidence spectrograph. No transverse spectrograph was used so that gain measurements were not possible. For a few shots the fibre was tilted at a small angle and the cylindrical lens rotated to form the line forms at this angle. Comparison of these "transverse" spectra with the axial spectra suggested a gain length product $g\ell \sim 1$, although the uncertainties in focusing and energy coupling make quantitative comparison difficult. Nevertheless, these results were encouraging.

For the four beam experiment two spectrographs would be used and field of view considerations meant that it was preferable to use a spectrograph with high sensitivity at $182\text{\AA}$ for the transverse position. This was because the transverse spectrograph had to be further from the plasma than the axial instrument. The single beam experiment allowed a qualitative comparison of the sensitivities which showed that the QUB spectrograph was better at $182\text{\AA}$ than the Hull instrument by a factor ~ 2 . Consequently, the QUB instrument was more suitable for the transverse position. The resolution of either instrument was expected to be insufficient to resolve the H α "doublet" but the higher dispersion of the Hull spectrograph would be more likely to show any structure which did appear in the axial spectrum. A number of shots in the single beam experiment did show a sharp central feature in the H α line which was better seen in the four beam experiments.

7.2.3(a) Experiment with multiple green beam illumination

This experiment was designed to measure gain on the C VI H α line ($182\text{\AA}$) in a symmetrically cylindrically expanding carbon fibre. Fibres of approximately $5\text{ }\mu\text{m}$ diameter were aligned at the common focus of four green beams which were focused into $\sim 3.3\text{ mm}$ line foci using a combination of a -3.3 m focal length plano-concave cylindrical lens and a 10.8 cm spherical lens in each beam (Fig. 7.03). In practice this focus was marked down so that only the central region of the gaussian beam was used, thus

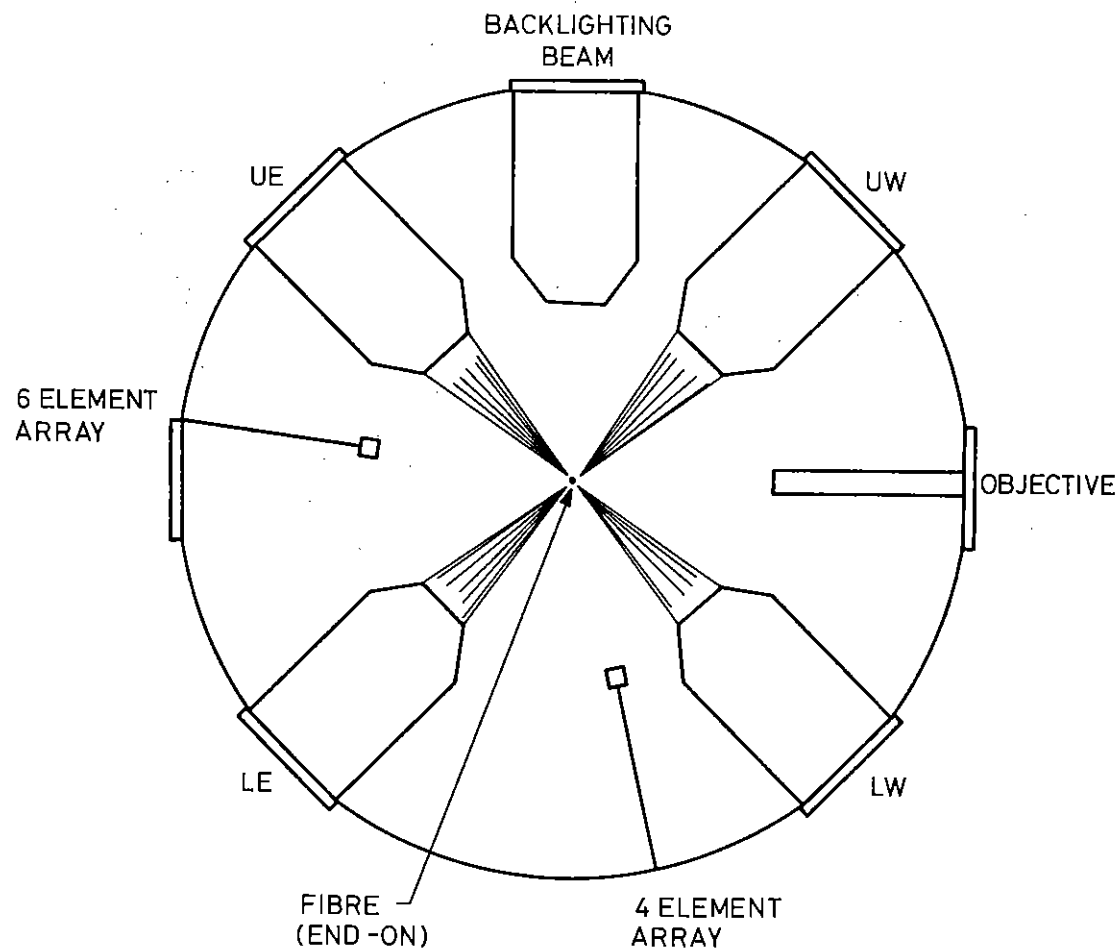


FIG. 7.03. VIEW FROM NORTH (LOOKING ALONG FIBRE AXIS)
(SPECTROGRAPHS OMITTED)

preventing the formation of cold absorbing plasma by the low intensity wings of the beam. The laser pulse duration was 200 ps, with a 20% prepulse. The delay time between pulses was 400 ps. Up to 10 J of second harmonic energy was available per beam with the beam energies balanced to better than $\approx 10\%$. The heated fibre length was ≈ 1.8 mm. Fibre alignment was greatly simplified with the use of visible green light from a cw argon-ion laser and the availability of a specially designed goniometer mount for supporting the fibre. A North-South axis, in which the fibre would lie, was initially defined by stretching a 50 μ m diameter wire along a target chamber diameter. The orientation of the 4 cylindrical lenses was adjusted to give total obscuration of green light in the near field using opposite lens pairs. The defining wire was imaged to a distant reference plane with a high resolution microscope objective at $\times 100$ magnification. The fibre was then positioned to first order accuracy using the reference position and fine tuned with θ , ϕ , x , y and z goniometer movements to simultaneously give the maximum obscuration with the highest symmetry from each lens. The free end of the fibre pointed directly into an axial spectrograph.

7.2.3(b) Gain Measurements

Gain measurements were made using two 1 m grazing incidence spectrographs arranged as shown in Fig. 7.04. The first looked axially along the length of the plasma, the second viewed the plasma transversely at an angle of 50° . Because of the high aspect ratio of the plasma at the onset of gain, amplification will occur along its length and will therefore be manifested by an anomalously high intensity H α signal recorded by the axial spectrograph. In practice we looked for an anomalous ratio of axial signal/transverse signal. The spectrographs were cross-calibrated using the backlighting beam in a spot focus on a suitably angled solid carbon target to establish what this ratio would be in a non-gain medium. In the cross-calibration experiments, precautions were taken to ensure that both spectrographs viewed exactly the same plasma.

If I_A^F , I_T^F and I_A^C , I_T^C are the axial and transverse signals for fibre shots

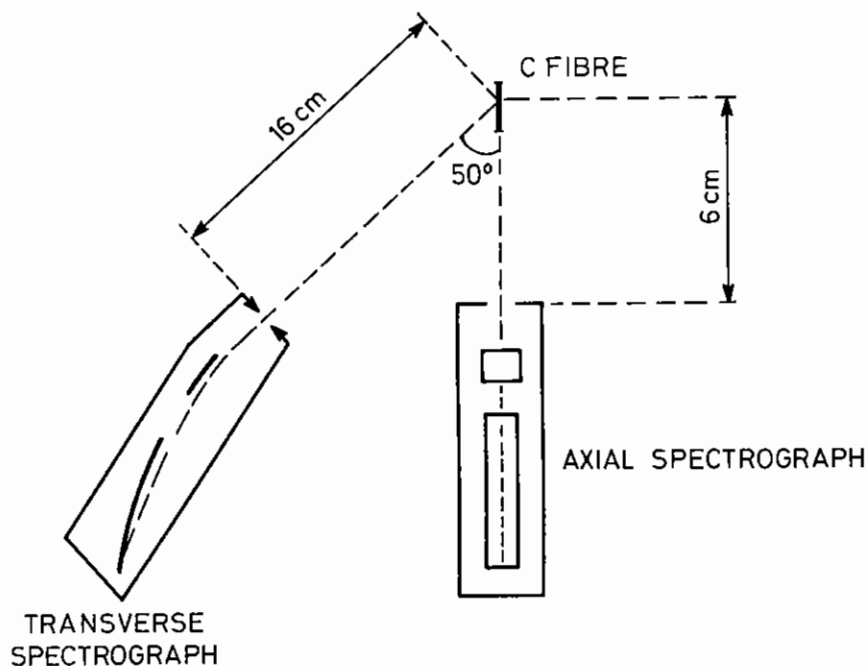


FIG. 7.04 . LAYOUT OF TRANSVERSE AND AXIAL SPECTROGRAPHS

and cross-calibration shots respectively, then we may show that the amplification is given by

$$G = \frac{I_A^F / I_T^F}{I_A^C / I_T^C} = \frac{e^{\alpha l} - 1}{\alpha l}$$

where α is the gain coefficient (cm^{-1}) and l is the length of the active medium (cm).

The energy coupled into the plasma was measured using the two plasma calorimeter arrays used in the previous single green beam illumination experiment. Their arrangement is shown in Fig. 7.03. The 6 element and 4 element calorimeters were respectively 20 cm and 23 cm from the target. The energies measured by the calorimeters indicated symmetrical, highly directional expansion with an average energy coupling of $> 10\%$, the maximum coupling was $\sim 24\%$ (Fig. 7.05). The plasma flow was contained within an angular spread of $< 10^\circ$ in the plane perpendicular to the fibre axis.

7.2.4 Data and Analysis

A computer program was developed at QUB and used with Joyce-Loebl auto-densitometer to integrate the intensity under the spectral lines of interest. Integration over line profiles to measure the relative energy emitted in a given line transition required a knowledge of the spectrograph film (Kodak 10101) characteristic response curve. This calibration was done using the QUB Nd/glass laser and will be reported elsewhere (7.03). The main conclusion of interest here was that the film responded linearly up to exposed densities of ≈ 0.5 D above film fog for radiation at $182 \text{ \AA}$. Since the data in the present experiment was within this linear response region relative line intensities were evaluated by calculating $\int D \, dx$ over a line profile. Figure 7.06 shows the integrated axial vs transverse H α signals. The solid line is that obtained from the cross-calibration points, the dotted lines represent the maximum errors in the cross-calibration ratio. The fibre shots are indicated by crosses,

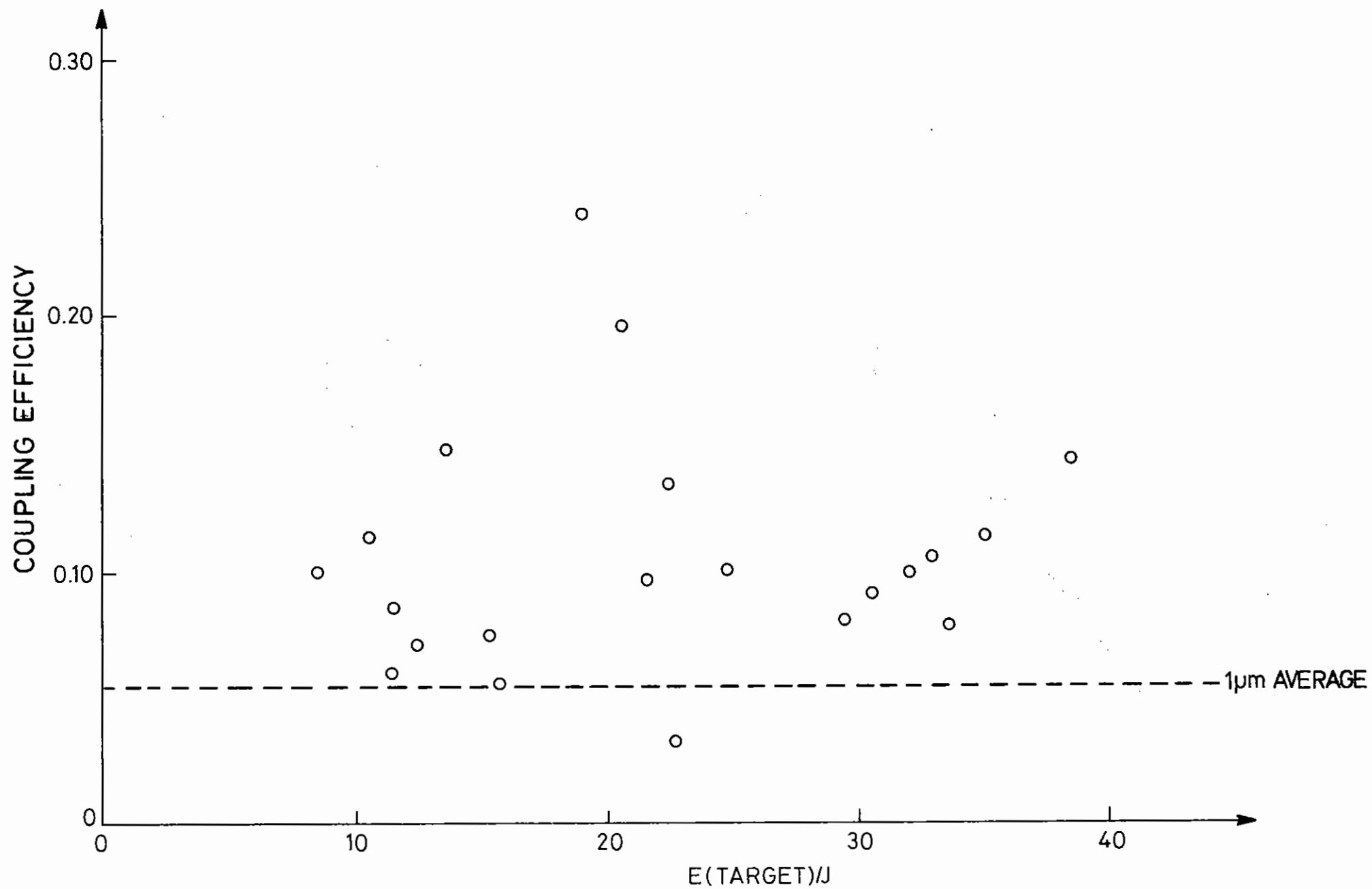


FIG.7.05. COUPLING EFFICIENCY AS A FUNCTION OF INCIDENT ENERGY ON TARGET

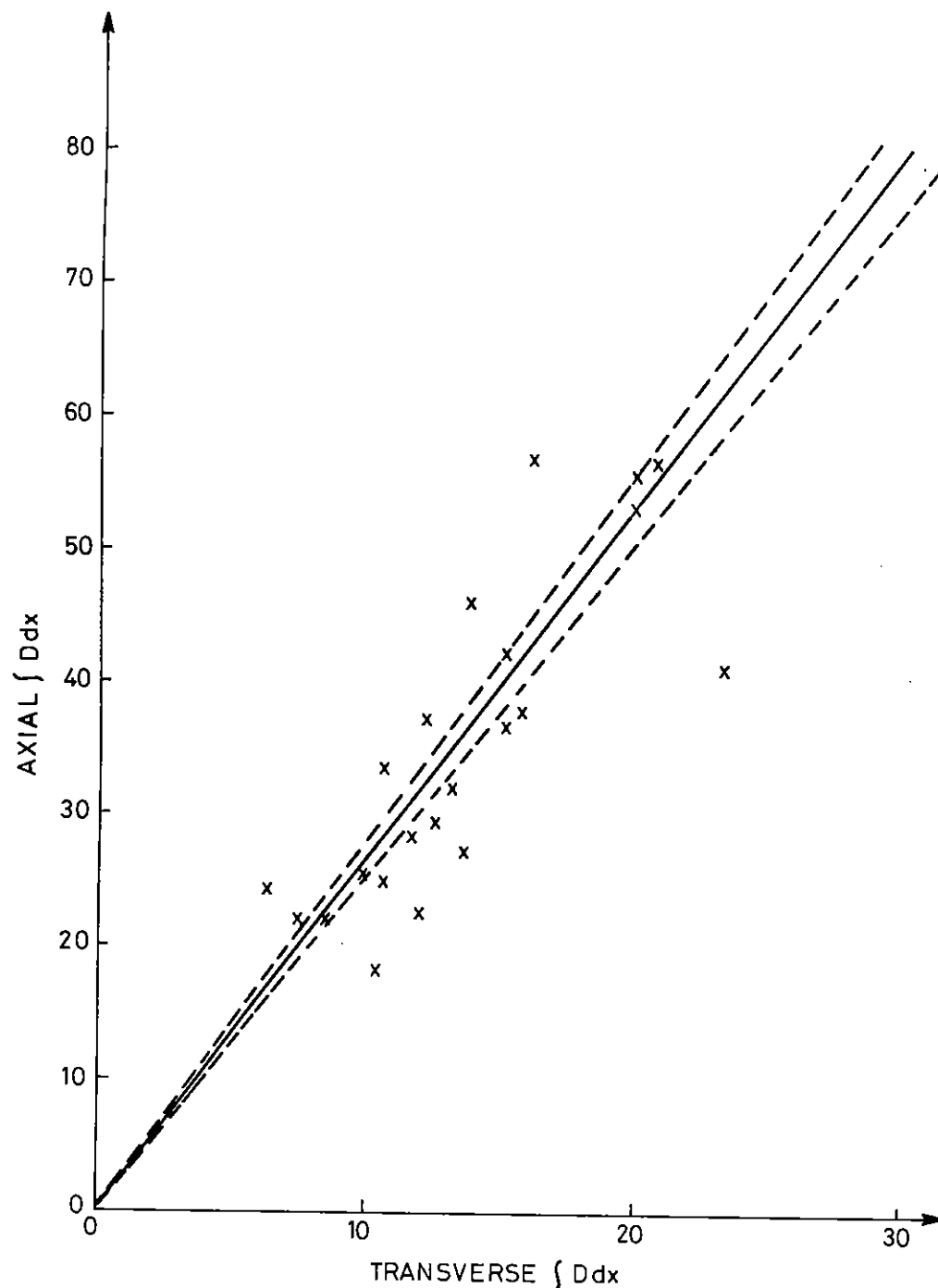


FIG. 7.06. INTEGRATED AXIAL vs INTEGRATED TRANSVERSE
H α SIGNALS

those lying above the cross-calibration line show amplification, those below show absorption. Fig. 7.07 summarizes the values of αl obtained, the largest being $\alpha l \sim 0.7$ ($G = 1.44$) at 0.5 J/mm energy absorbed. Absorption is chiefly due to non-uniformity of plasma temperature along its length as shown in the pinhole picture of a heated fibre, (Fig. 7.08). The colder plasma lying in the gain path absorbs the radiation and destroys any gain. The presence of this cold plasma is discussed in more detail in section 7.2.5.

The irreproducible behaviour of the gain and its apparent independence on absorbed energy is basically due to poor design of the experiment. The fibre diameters used were not properly matched to the laser energy available. Using too thick fibres meant that a relatively large absorbing cold residual core was left after the expansion, whose behaviour was erratic and unpredictable (cf Section 7.2.5).

A very interesting observation is that in the fibre shots which did exhibit gain, there was a central feature (resolution limited) on the H α line (Fig. 7.09). This structure was present only on the H α line and only on the axial spectra. The fact that it was not observed on the Ly β line which is also optically thin would support the suggestion that it is attributable to gain on the H α line. Close examination of the axial spectra obtained in the previous single green beam experiment also showed this central line structure, although gain measurements were not made on this occasion.

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7.2.5 Computer Modelling of Fibre Under Green Beam Illumination

In order to model the four beam fibre heating experiments a series of simulations using the codes POLLUX and GAIN have been carried out. The two dimensional hydro code POLLUX was modified to consider one quadrant of the cylindrical fibre only. Although the code considers only the hydrodynamics of a two fluid ideal gas, a simple correction for ionisation

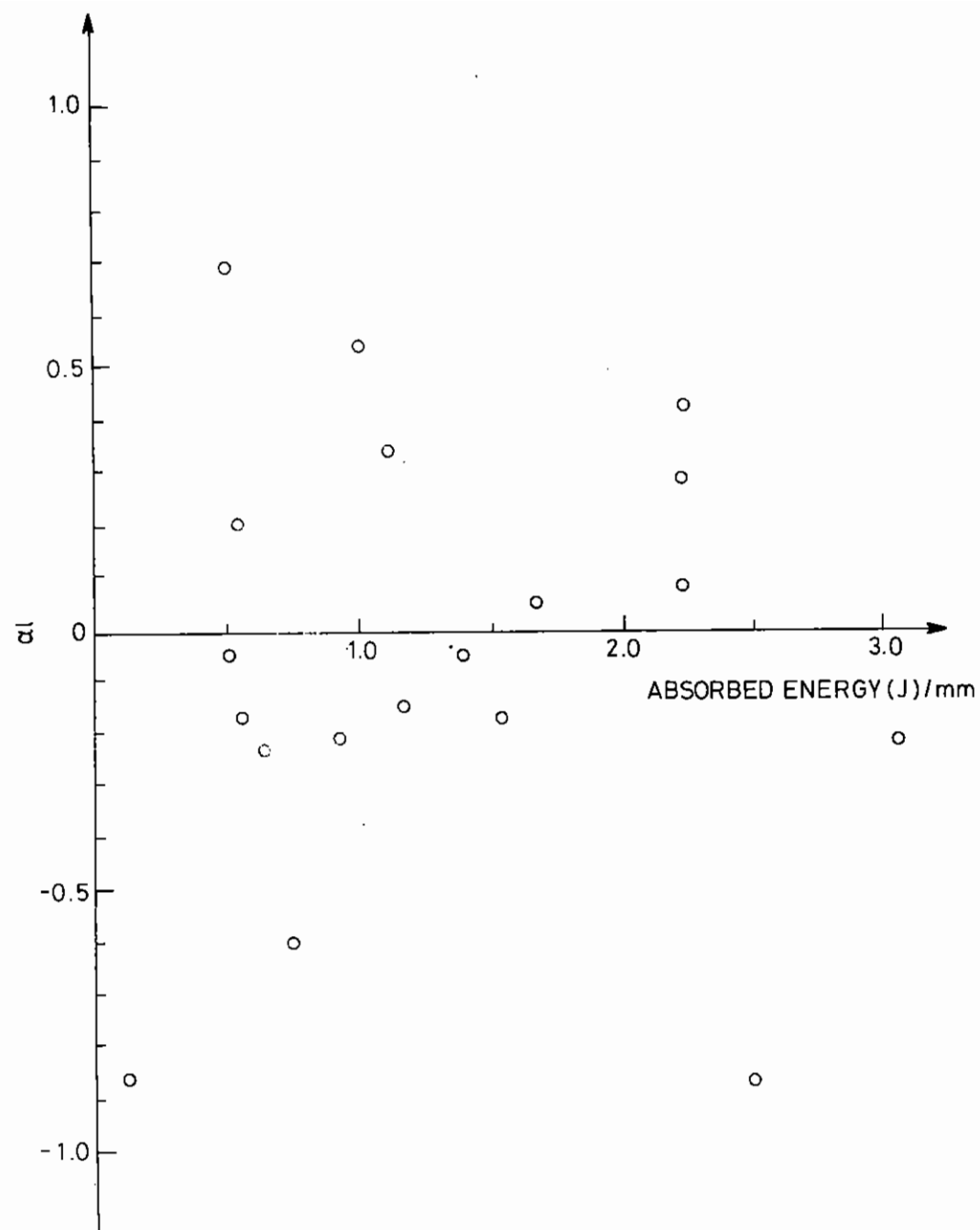


FIG. 7.07. VALUES OF αI OBTAINED AS A FUNCTION OF THE ABSORBED ENERGY

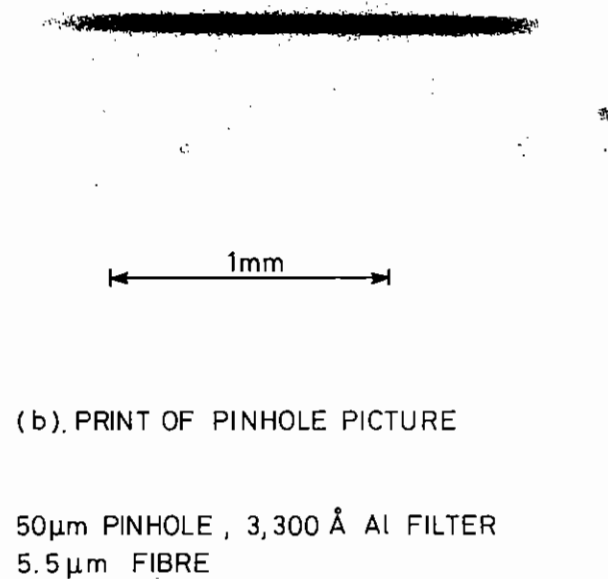
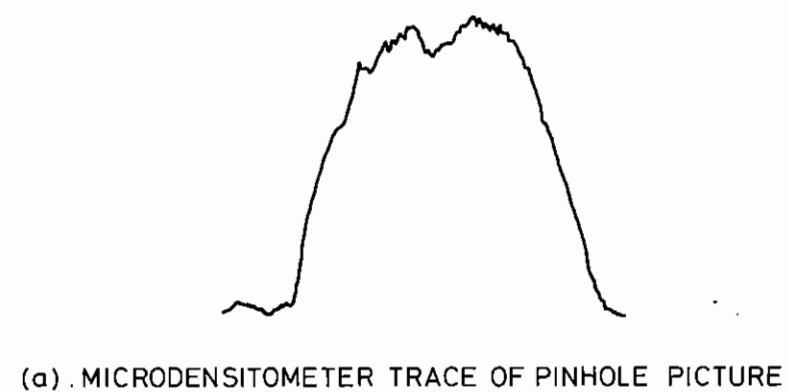


FIG. 7.08 . X-RAY PINHOLE PICTURE OF HEATED FIBRE

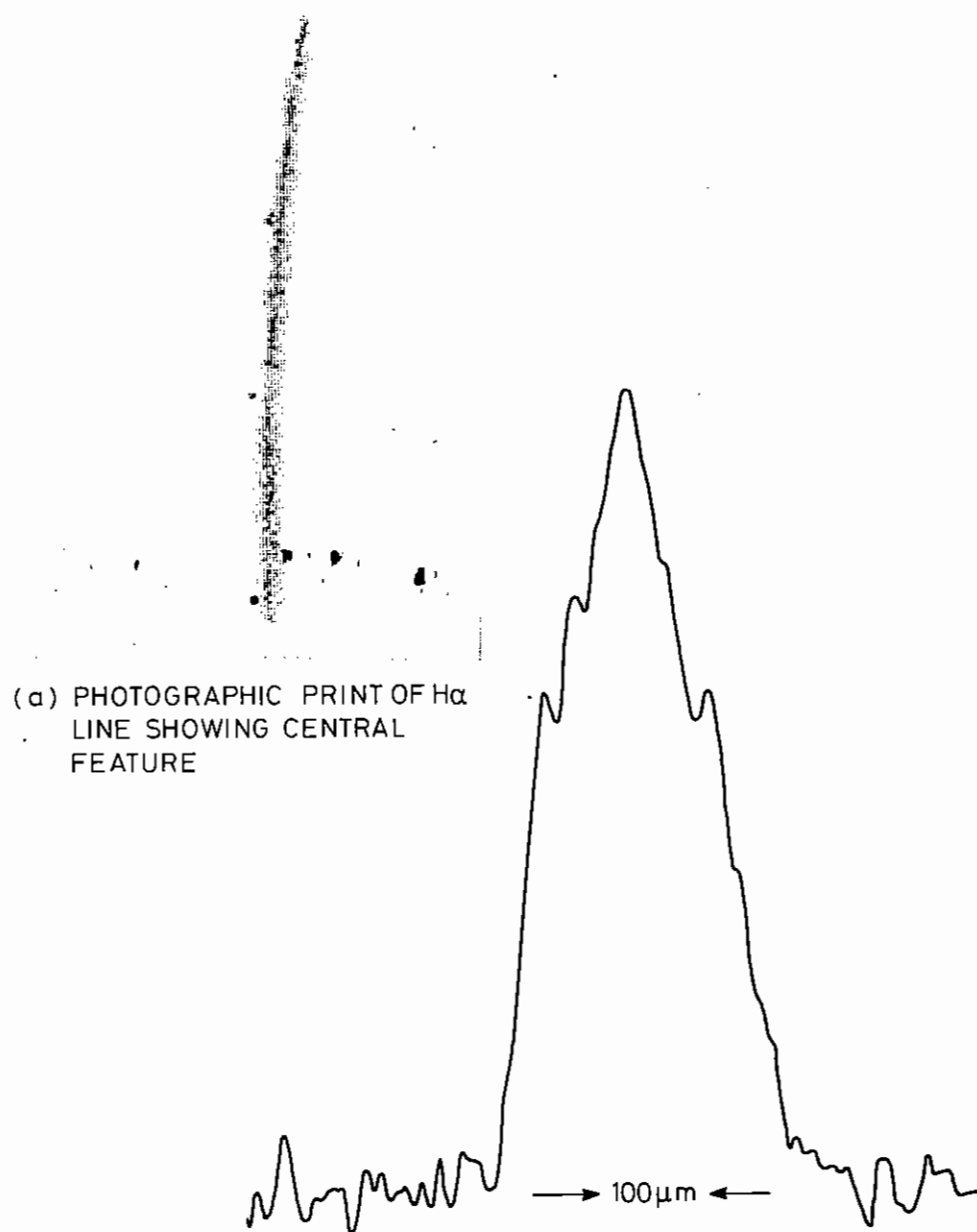


FIG. 7.09. CENTRAL FEATURE ON H α LINE

could be made to allow the data to be compared with the results from the ionisation code GAIN.

The basic behaviour of the flow in a fibre heated by laser was described in reference (7.01). In a thick fibre the burn is incomplete, and at the end of the laser heating a cold dense core (residual core) is left surrounded by a blanket of hot plasma. Once the heating phase ceases, the hot plasma expands, and heating of the residual core by thermal conduction ceases. Since the residual core is relatively small it plays little part in the subsequent expansion of the hot plasma, which if well formed is nearly isothermal and symmetric. In this case the expansion of the hot plasma is well described by the similarity model used in the code GAIN. The development of ionisation may thus be calculated with the latter code provided the hot plasma mass - effective mass - is used instead of the actual fibre mass; an hypothesis justified by experiment (7.01).

In Figure 7.10 we present results obtained in this way. The effective mass m_{eff} calculated by POLLUX, and corrected for ionisation m_{eff}^i are plotted as functions of the absorbed laser energy per cm of fibre length. This may be compared with the graph of the required absorbed energy, G_{max} , for maximum gain. The data is calculated for a 5 μm diameter fibre irradiated by a 0.53 μm laser of pulse length 200 ps, with a 20% pre-pulse of delay 400 ps, corresponding to the actual experimental conditions used in the four beam experiments. It is immediately apparent that there are two regimes of operation - one at high absorbed energy ($\sim 40 \text{ J/cm}$) in which the fibre is completely burnt through, and one at low ($< 5 \text{ J/cm}$) in which only a small fraction of the fibre is heated.

In order to assess the importance of non-optimisation we have also plotted the required energy to give a gain peak half the optimum, a value which in view of the nature of the gain profile, corresponds roughly to the limit of detectable gain in a time integrated experiment. It can be clearly seen that in the range of absorbed energy 10 - 30 J/cm no gain should be detected.

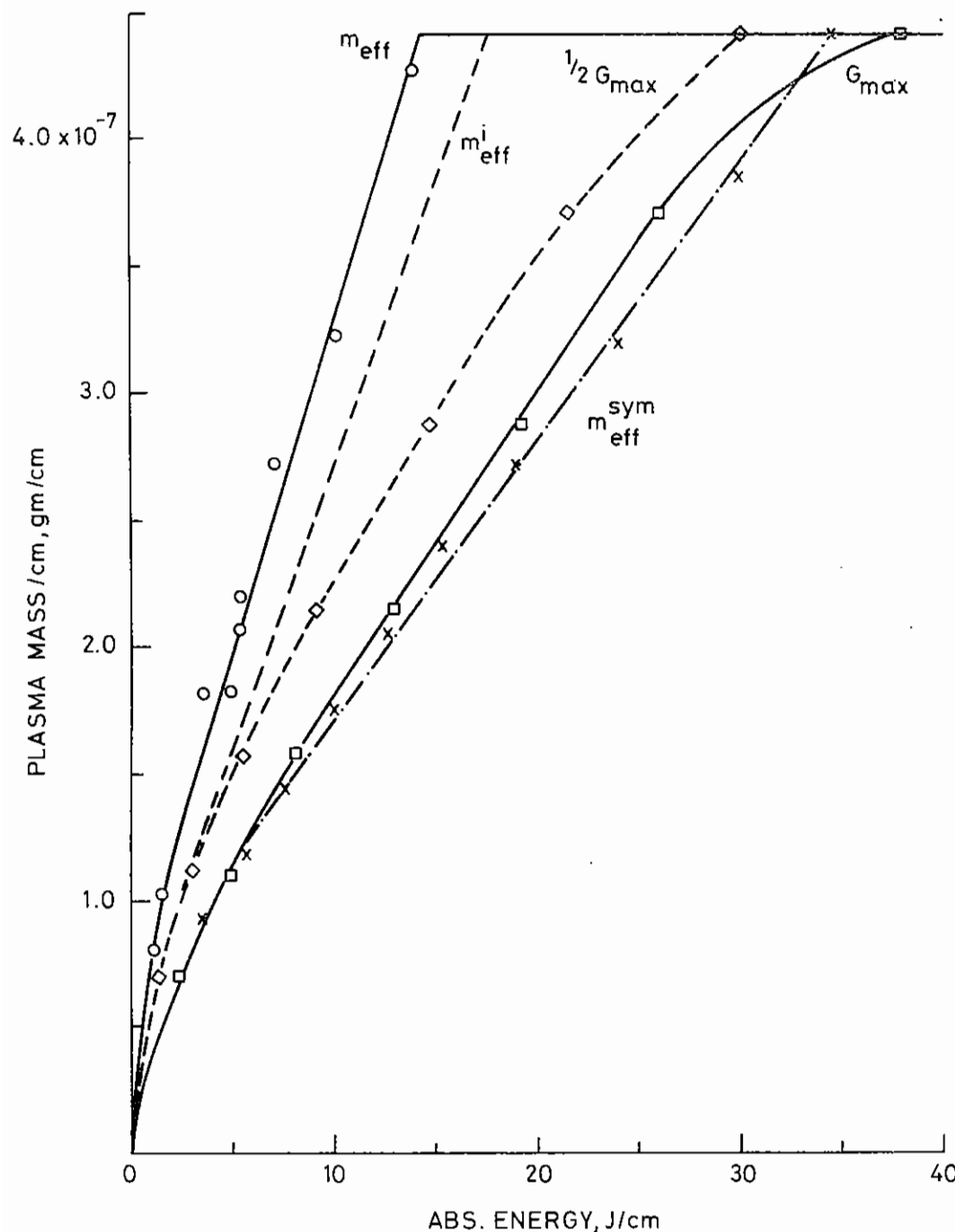


FIG. 7.10. EFFECTIVE MASS AS A FUNCTION OF ABSORBED LASER ENERGY PER cm OF FIBRE LENGTH. $0.53 \mu m$ BEAM ILLUMINATION

If the fibre could be heated symmetrically these results could be markedly changed, as shown by the effective mass, m_{eff}^{sym} , under symmetric heating. In this case the heated mass curve closely matches the required energy curve, and good gain should be achieved over the entire range. The difference in these two modes of heating appears to be due to eddy structures developed by the higher total intensity on the diagonals of the illumination pattern: checks show it is not a code effect.

The behaviour described above is quite general, thus Figure 7.11 shows an almost identical heating pattern for a single beam one sided illumination of $1.06 \mu m$ radiation, pulse length 150 ps, with a 20% pre-pulse, delay 200 ps. In this case also gain is detectable only if the absorbed energy is less than 10 J/cm or greater than 35 J/cm. This system corresponds to the Hull experiments (7.01), which were performed with thinner fibres. Here the low energy regime corresponds to non-symmetric heating and therefore to an unreliable condition due to the behaviour of residual core.

The optimum gains obtained on the matched heating curves, G_{max} , are plotted in Figure 7.12 for the above cases. The difference in gain between the two situations is primarily due to the longer pre-pulse in the Central Laser Facility experiments, which causes the gain to occur later, and therefore at lower density, in the expansion.

In a practical system at high gain the residual core, which is cold and therefore strongly absorbing, may seriously interfere with the optical path for laser action. If the beam intensities (in a multiple beam experiment) are not well matched along the entire length of the fibre, regions of cold core may be driven into the gain path and prevent laser action. Although these effects are avoided by one sided illumination other considerations probably rule out this solution. The problem may obviously be eliminated if the fibre is burnt through. However the use of thinner fibres, such that thermal conduction enhances the symmetry of the burn, should reduce core movement into the gain region, and overcome this problem.

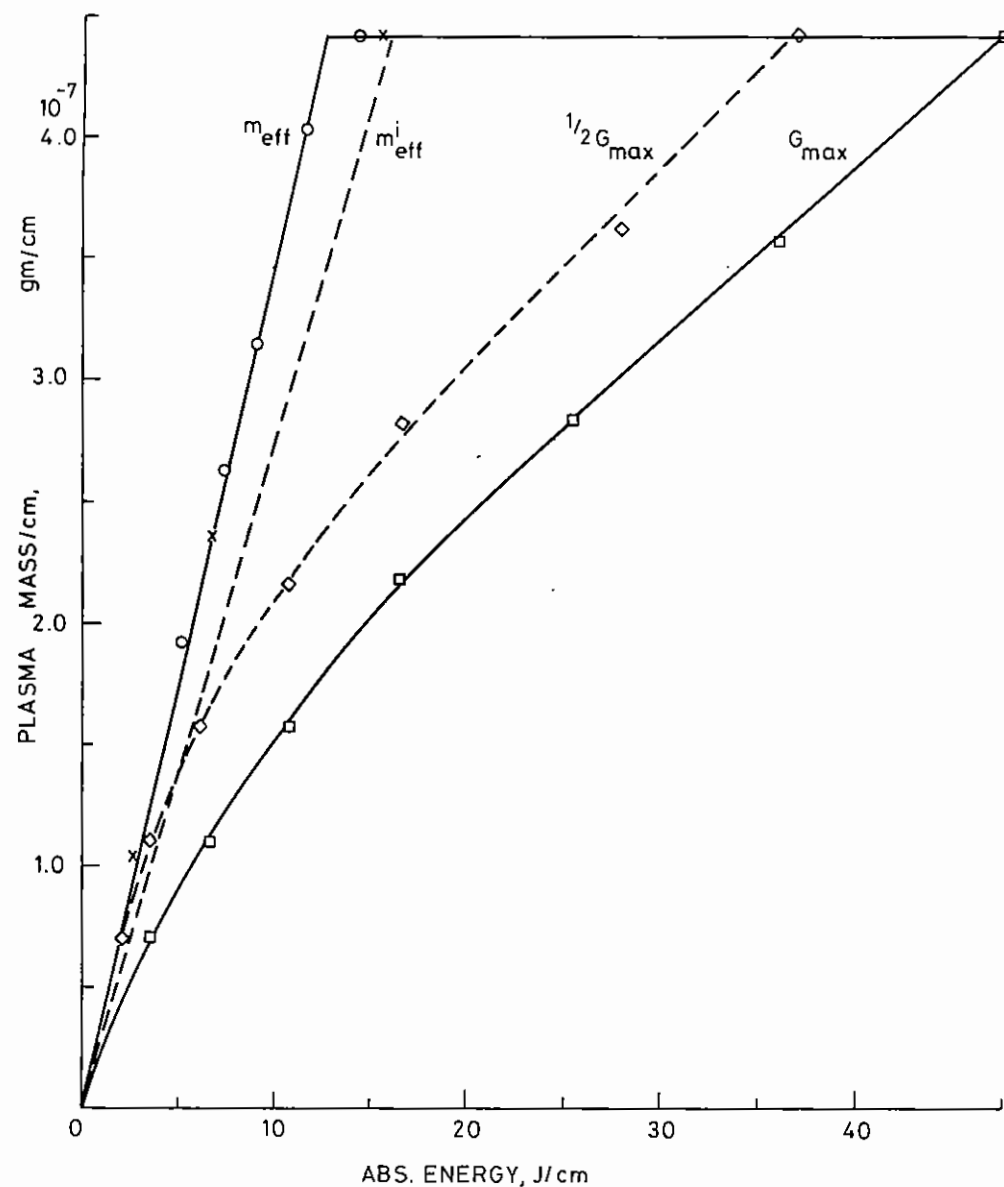


FIG. 7.11 EFFECTIVE MASS AS A FUNCTION OF ABSORBED LASER ENERGY PER cm OF FIBRE LENGTH. $1.06 \mu\text{m}$ BEAM ILLUMINATION

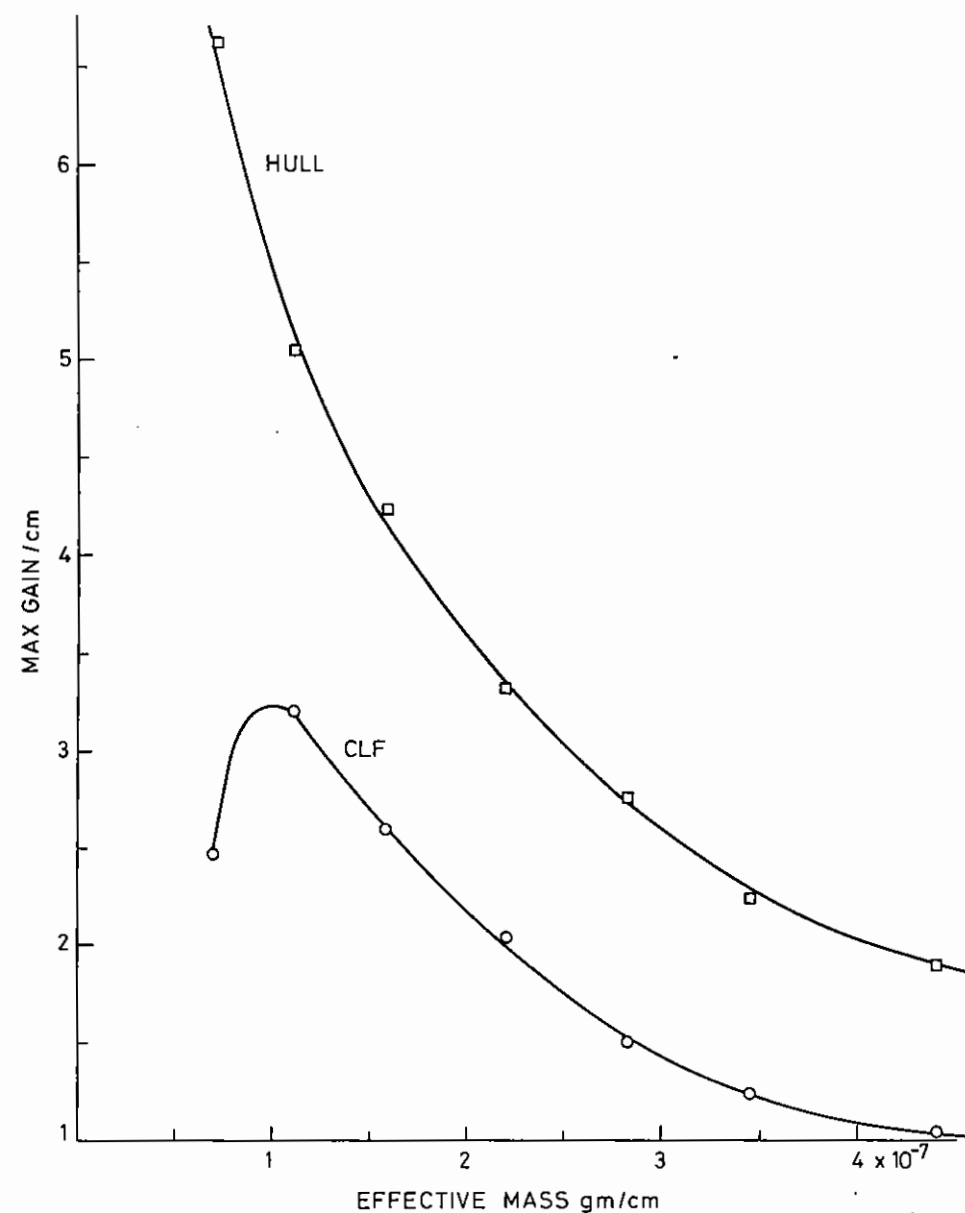


FIG. 7.12. OPTIMUM GAINS OBTAINED ON MATCHED HEATING CURVES FOR HULL AND CLF EXPERIMENTS

The behaviour observed in the four beam experiments is in good accord with these simulations. Gain was observed, but irreproducibly and of low value in the two spectrograph experiments, occurring principally at the lower absorbed energies 5-10 J/cm. In this context the observation of a narrow (resolution limited) line superimposed on a broad emission at H_{α} occurring only on H_{α} and in the axial spectrograph alone is important. The broad H_{α} emission must occur earlier in the expansion, and corresponds to a plasma inadequately heated for its mass. However even in such cases gain will set in (but may not be measurable in a comparative test) at a later, low density cold stage of the expansion. This effect probably accounts for the narrow line, which exhibits gain with an (αk) of about 2.

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7.3 Dielectronic Satellite Spectra

7.3.1 Density-dependent Si XII dielectronic satellite spectra

In the laser implosion of gas-filled glass microballoon targets both the ablation plasma and the implosion core plasma are intense sources of x-ray line radiation from multiply-charged ions. The radiative decay of doubly-excited levels of helium-like ions results in the emission of spectral lines that are long-wavelength satellite lines of the hydrogenic resonance lines. It has been shown (7.04) that the relative intensities of the Si XIII satellite features of the Si XIV Lyman- α line vary with electron density in the plasma, and represent a promising density diagnostic technique for electron densities in the range 10^{22} to 10^{24} cm⁻³. It has also been predicted (7.05) that the relative intensities of the Si XII satellite spectrum on the long-wavelength side of the $1s2p\ ^1P-1s^21S$ resonance line of the helium-like ion Si XIII are density dependent. This letter reports the experimental observation of these density-dependent effects in the x-ray spectrum of Si XII emitted from laser-imploded microballoons.

The relative intensities of the satellites of the helium-like $1s2p\ ^1P-1s^21S$ line have been calculated for medium-Z elements by Gabriel (7.06) and Bhatta et al (7.07) using a corona model theory, which assumes that the upper levels in the radiative transitions are populated only by dielectronic capture and inner-shell electron collisional excitation. This calculation leads to relative satellite intensities that are independent of electron density. In a more complete calculation (7.05) electron collisional transitions between closely-lying autoionising levels are also taken into account. For electron densities exceeding 10^{22} cm⁻³ collisional transitions between the autoionising states lead to an enhancement of the intensities of the $1s2p^2\ ^2P-1s^22p\ ^2P$ and the $1s2p(^1P)2s^2P-1s^22s\ ^2S$ satellite lines. These are denoted (a-d) and (q,r) respectively, using the key letter description of Gabriel (7.06). It is also predicted that for electron densities up to 10^{24} cm⁻³ the intensity of the $1s2p^2D-1s^22p\ ^2P$ lines (denoted by k, l, j) are only weakly affected by these collisional transitions between the autoionising levels, and that the sum of their intensities relative to the $1s2p\ ^1P-1s^21S$ helium-like resonance line is mainly a function of the electron temperature. Thus, for electron densities in the range 10^{22} - 10^{24} cm⁻³ it is predicted that the intensities of the (q,r) and (a-d) satellite lines will increase relative to the (k,l,j) satellites.

The Si XIII $1s2p\ ^1P-1s^21S$ and associated satellite line intensity distributions were obtained from the space-resolved x-ray spectra recorded for a series of two-beam laser implsoions at the Rutherford Laboratory (7.08). The parameters of three representative laser shots are shown in table 7.01. The electron temperature is calculated from the slopes of the Si XIII and Si XIV recombination continua, and the electron density is determined from the Stark broadening of the Si XIII and Si XIV resonance lines, corrected for optical opacity (7.08).

The $1s2p\ ^1P-1s^21S$ and satellite line spectra from the ablation plasma on shot 328 and the implosion core plasmas on shots 323 and 321 are shown in figure 7.13. It is clearly seen that the intensities of the (q,r) and (a-d) satellites relative to the (k,l,j) satellites increase with increasing electron density. This effect is reproducible, as is observed in the

Table 7.01 Parameters for the three laser shots.

Shot number	Laser power (TW)	Shell diameter (μm)	Shell thickness (μm)	Gas fill (bar)	Plasma parameters	
					Electron temperature (eV)	Electron density (cm^{-3})
328	0.34	67	0.71	0.4 Ar	470	$1.2 \times 10^{22+}$
323	0.17	50	0.59	10 DT	580	$2 \times 10^{22\#}$
321	0.19	51	0.53	10 DT	510	$8 \times 10^{22\#}$

a The shot numbers refer to the numbering system for laser shots used at the Rutherford Laboratory.

+ Ablation plasma.

Implosion core plasma.

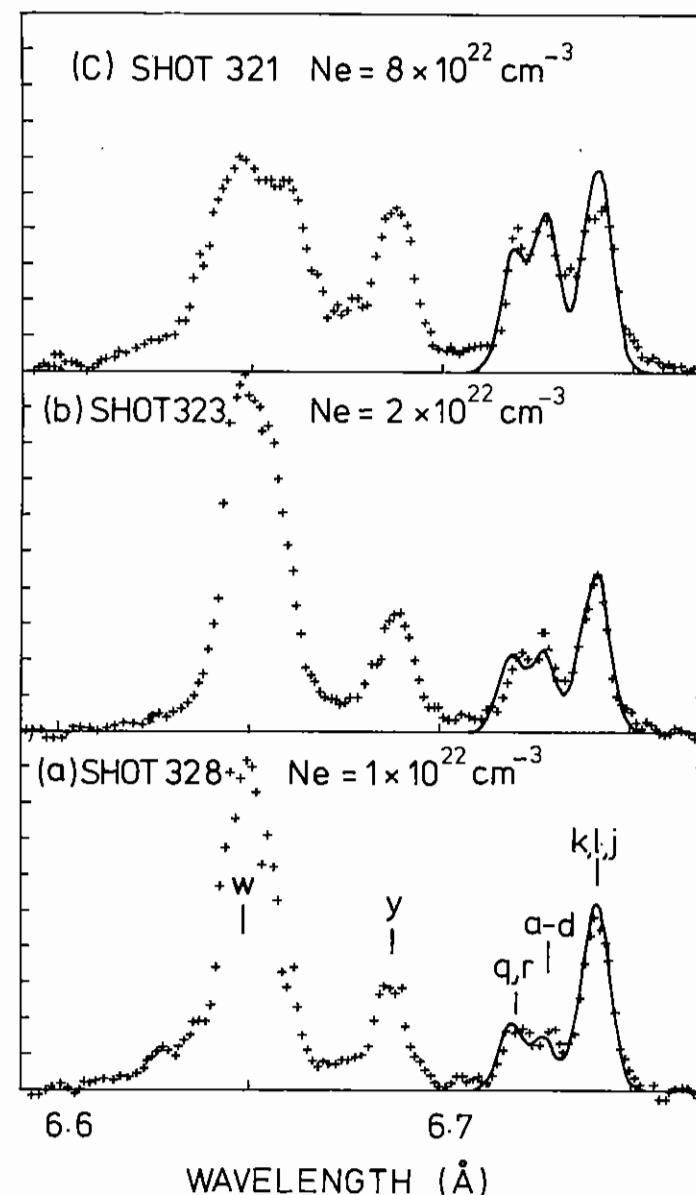


Fig. 7.13 Experimental spectra showing the Si XIII $1s2p\ ^1P-1s2^1S(w)$, intercombination (y), and dielectronic and inner-shell satellite lines for the three shots shown in table 7.01. Spectra (a) and (b) are from the implosion core plasmas in shots 321 and 323 respectively, while spectrum (c) is from the ablation plasma on shot 328. The full curves show theoretical fits to the satellite features as described in the text.

spectra recorded from six other implosion core plasmas and a much larger number of ablation plasmas. The spectra shown in figure 1 are representative of the electron density range over which the observations extend.

To compare the relative intensities of the satellite lines with theory a theoretical spectrum is computed at each electron density and fitted to the experimental spectra. The generalised emission rate coefficients describing dielectronic recombination, inner-shell electron collisional

excitation, and collisional mixing between the autoionising levels as a function of electron temperature and density are tabulated for the spectral lines of interest here (7.05). This tabulation is used to calculate the relative satellite intensities as a function of electron density for an electron temperature of 500 eV. Since the temperature dependence of the satellite line ratios is relatively weak a calculation at a single temperature is adequate for comparison with the three spectra shown in figure 7.13. In the calculation of the theoretical spectrum the wavelengths of the satellite lines are taken from Vainshtein and Safronova (7.09). In order to fit the linewidths of the experimental spectra the theoretical spectrum was computed using a Gaussian profile with FWHM equal to $7 \times 10^{-3} \text{\AA}$ on each line. The instrumental linewidth, due to the crystal rocking curve and the source size, is only about $5 \times 10^{-3} \text{\AA}$, which suggests that these satellite lines may be Stark broadened at the electron densities encountered here. The experimental spectrum is fitted by normalising the theoretical spectrum on the (k,l,j) spectral feature and then varying the fraction of Li-like silicon to fit the spectral feature comprised of the satellites (q,r) and (a-d). Inner-shell excitation makes a small contribution to the (k,l,j) satellites, but is more important for the (q,r) and (a-d) satellites. On shot 328 the population density ratio is

$$(n(1s^2 2s^2 S) + n(1s^2 2p^2 P)) / (1s^{21} S) = 0.03,$$

while on shot 323 and 321 this ratio is 0.065. At the electron densities encountered here, electron collisions will ensure that the $1s^2 2s^2 S$ and

$1s^2 2p^2 P$ states are populated in proportion to their statistical weights. In a calculation of the steady-state ionisation balance in a silicon plasma with an electron temperature of 500 eV and an electron density of 10^{22}cm^{-3} it is found that the population density ratio is (7.10)

$$(n(1s^2 2s^2 S) + n(1s^2 2p^2 P)) / (1s^{21} S) = 0.01$$

That calculation uses the ionisation model described by Stravrakas and Lee (7.11). The apparently high fraction of Li-like silicon revealed by the experimental spectra is probably due to the inhomogeneous nature of the plasma; there being a region of cooler plasma with more Li-like and less He-like and H-like ions than the region of plasma that is diagnosed by the He-like and H-like line emissions. Figure 7.14 shows the experimental values of the intensity of the sum of the (q,r) and (a-d) satellites relative to the (k,l,j) satellites as a function of electron density. The theoretical ratio calculated from (7.05) is also shown for two cases. In Figure 7.14 curve A shows the satellite ratio as a function of electron density when the number density of Li-like silicon is taken to be zero, so that inner-shell collisional excitation makes no contribution to the satellite intensities. Clearly, such a calculation does not agree with the experimental observations. Also in Figure 7.14, curve B shows the satellite intensity ratio when inner-shell collisional excitation is included by taking

$$(n(1s^2 2s^2 S) + n(1s^2 2p^2 P)) / (1s^{21} S) = 0.065$$

There is good agreement between theory and experiment for the increase in the satellite intensity ratio

$$(I_{q,r} + I_{a-d}) / I_{k,l,j}$$

with increasing electron density. This increase is due to electron collisional transitions between the autoionising states. In the experimental spectra the increase of the (a-d) satellite group relative to the (q,r) satellite group is less than is predicted by theory. This may be due to optical opacity on the satellite transitions, in a manner

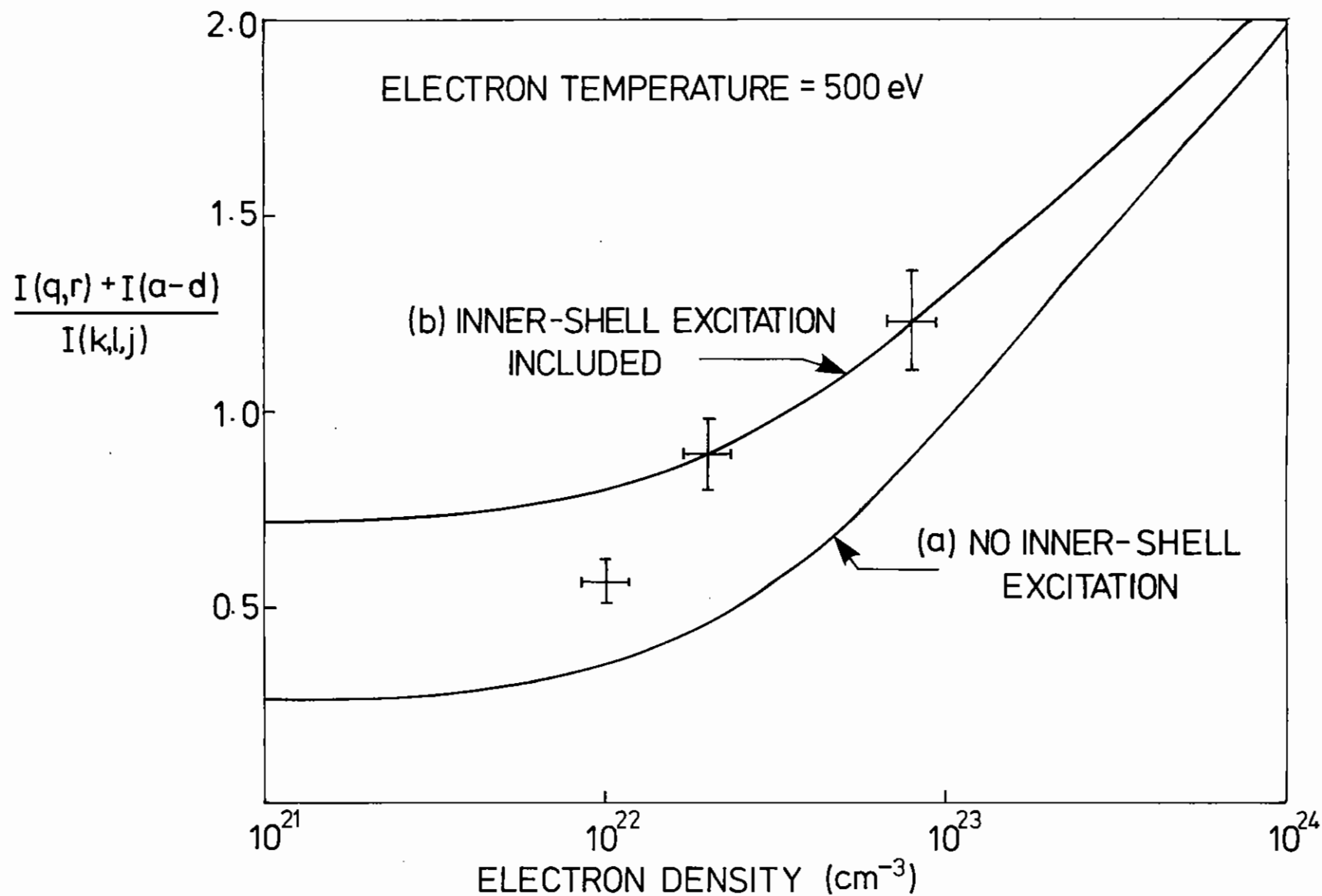


Fig. 7.14 The intensity ratio of the sum of the (q,r) and (a-d) satellite lines relative to the (k,l,j) satellite at an electron temperature of 500eV for two cases: A, without inner-shell excitation; B, with inner-shell excitation using $(n(1s^2 2s^2 S) + n(1s^2 2p^2 P)) / n(1s^2 1s^1 S) = 0.065$. Also shown are the experimentally observed intensity ratios.

analogous to the effects of optical opacity on the Lyman- α satellites (7.12). Considering only Doppler broadening it is estimated, for example, that for shot 323 the peak opacity of the (a) satellite is about one. For plasma density diagnosis it is clearly preferable to arrange for the satellite lines to be optically thin. The usefulness of these satellite line intensity ratios for the diagnosis of the plasma density is somewhat diminished by the fact that the satellite intensity ratio depends on both the electron density and the population density ratio between helium-like and lithium-like ions.

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7.3.2 The Density Dependence of Dielectronic Satellites of Helium-Like Transitions

7.3.2(a) Introduction

Dielectronic satellites to helium-like resonance lines are a prominent feature in both low (7.13) and in high density (7.14) plasmas. The transitions arise from three electron spectra of the form $1s^2n\ell - 1s2\ell' n\ell$. When the spectator electron $n\ell$ has principal quantum number $n = 2$, there are about 26 lines. At low density their relative population is essentially coronal (7.15) whereas at high density the levels become statistically populated. There is, however, an intermediate region where the population ratio of groups of dielectronic satellites is a function of the electron density (7.05). Moreover, because of the low opacity of the dielectronic satellites, their intensity can be used as a diagnostic of the level population. A similar dependence has been observed for the satellites to the hydrogen like resonance lines (7.16).

We present for the first time results of new calculations of the rate equations for arbitrary Z plasmas. The excited levels of Li, He and H like ions are represented. The most accurate detailed rates available are used and results are presented for the relative intensities of Li like satellite to the He like 1-2 transition for a variety of materials useful for laser plasma diagnosis. It is shown that useful density diagnosis can

be achieved over a range of about 20 for any element if the resolution is sufficient to resolve emission from the abcd ($1s^22p^2 - 1s2p^2 2p$) complex from qr. ($1s^22s^2S - 1s2s2p(^1P)1s^2P$) (see Gabriel (7.06) for the nomenclature). Otherwise the ratio of the combined abcd + qr to jk ℓ intensity ($1s^22p^2 2p - 1s2p^2 2D$) gives a useful range of about 10 in the diagnosis of plasma density. The density at which the ratios are usefully density-dependent is about $3 \times 10^{21} \text{ cm}^{-3}$ for Na up to $4 \times 10^{23} \text{ cm}^{-3}$ for Ar. Such densities are readily achieved in laser plasmas.

The model is also used to predict the intensity of the dielectronic satellites lying close to the intercombination line ($1s^2 1p - 1s 2p^3P$). The results show that these underlying satellites make it difficult to measure the intensity of the intercombination line at high density. Results from the calculation are compared with spectra from ablation plasmas and the cores of exploding pushers.

Ablation plasmas produced by the interaction of 1 ns, 0.53 μm and 1.05 μm light at $10^{15} \text{ W cm}^{-2}$ with solids, emit X-ray at approximately the critical density, ie $4 \times 10^{21} \text{ cm}^{-3}$ and $1 \times 10^{21} \text{ cm}^{-3}$ respectively. The long time-scale of these plasmas allows an equilibrium ionisation model to be used with some confidence. A suitable element illustrating applicability of the model to the satellite line ratios is Na. In contrast, results from a Si spectra produced by an ablation plasma show that very good resolution is required to enable the low density to be diagnosed.

Higher densities can be achieved using so-called exploding pusher targets (7.17). Although existing for only 0.1 ns, these plasmas are at a high enough density to equilibrate the populations of different stages. Densities of $2.5 \times 10^{23} \text{ cm}^{-3}$ are produced by 1.05 μm irradiation, and densities up to $4 \times 10^{23} \text{ cm}^{-3}$ are produced with shorter wavelength irradiation. Experimental results are presented illustrating the satellite ratios for Cl and Ar, and are contrasted with a much lower density Cl ablation plasma. Again agreement with the theoretical model is demonstrated.

7.3.2(b) The Model

The rate equation model used for the ratio of the Li-like dielectronic satellites is based on a steady state equilibrium model for the ionisation balance. The atomic data required for the calculation are the energy levels, the collision cross-sections and the transition rates.

The model is solved for the ground states of all states of ionisation stages from the neutral atom up to the Be-like ion, and then several energy levels of the Li, He and H-like ions, and bare nuclei. The energy levels considered and their sources are shown in Table 7.02.

The collision rates are crucially important because their value determines the temperature and density at which the ratios of dielectronic satellites start to respond as density diagnostics. The sources of collisional rates between the different states are shown in Table 7.03. The most critical rates are between the autoionising levels of the lithium-like species. For these, the calculated cross-sections of Jacobs and Blaha (7.05) supplemented by the high energy effective Bethe Gaunt factor, are used. A parameterisation of cross sections, similar to that used by Sampson et al (7.25) is used to facilitate the integration over the Maxwellian velocity distribution. Since data are only available for Lithium-like ions of Al, the cross sections are scaled as Z_{eff}^2 , defined by Golden et al (7.24). This should not cause a substantial error over the range of Z considered here, ie Na to Ar.

7.3.2(c) Numerical Results

Figure 7.15 shows results of the calculations for Si. The ratios of the emission intensities of the three satellite complexes, abce, qr and mn, to the intensity of the dominant jkl complex, are shown as a function of density at a temperature of 500 eV. For the qr/jkl and abcd/jkl ratio results are also given at a temperature of 750 eV showing that the variation with temperature is negligible compared to the density variation. The mn/jkl ratio is also shown although the sensitivity to density is small.

TABLE 7.02

THE LEVELS INCLUDED IN THE MODEL

Ion Stage	Spectroscopic Specification	Comment	Source for Energies
Ions up to Li-like	Ground states		Kelly and Harrison (7.18)
Li-like	$1s^2 2s^2 S$	Ground state	
	$1s^2 2p^2 P$	Excited state	Clark et al. (7.19)
	$1s 2s^2^2 S$	Mixed, auto-ionising (op)	Gabriel (7.06)
	$1s 2p(^1P) 2s^2^2 P$	Mixed, auto-ionising (qr)	" "
	$1s 2p(^3P) 2s^2^2 P$	Mixed, auto-ionising (st)	" "
	$1s 2p^2^2 D$	Autoionising(jkl)	" "
	$1s 2p^2^2 P$	" (abcd)	" "
	$1s 2p^2^2 S$	Mixed, auto-ionising (mn)	" "
He-like	$1s^2^1 S_0$	Ground state	
	$1s 2s^3 S$		Extrapolation
	$1s 2s^1 S$	Dipole forbidden	Clark et al. (7.19)
	$1s 2p^3 P$	Intercombination	Scofield (7.20)
	$1s 2p^1 P$	Resonance	" "
	$1s 3$	Lumped state	Scofield (7.20)
	$1s 4$	" "	" "
	$1s 5$ (etc.)	" "	" "
	$2s^2^1 S$	Autoionising	Vainshtein & Safrannova (7.09)
	$2s 2p^3 P$	"	
	$2p^2^3 P$	"	
	$2p^2^1 D$	"	
	$2p^2^1 S$	"	
	$2s 2p^1 P$	"	
	$2s^2^1 S$	"	
Hydrogenic n =1 to 5		Lumped	Scofield (7.20)

TABLE 7.03

COLLISION RATES REFERENCES

TRANSITION	REFERENCE
Li-like bound to bound	Clark et al. (7.19)
Li-like bound to autoionising	Sampson et al. (7.21)
Li-like autoionising to autoionising	Jacobs and Blaha (7.05)
He-like bound to bound	Vinogradov et al. (7.22)
He-like bound to autoionising	Sampson et al. (7.21)
He-like autoionising to autoionising	Clark et al. (7.23)
Hydrogenic bound to bound	Golden et al. (7.24)

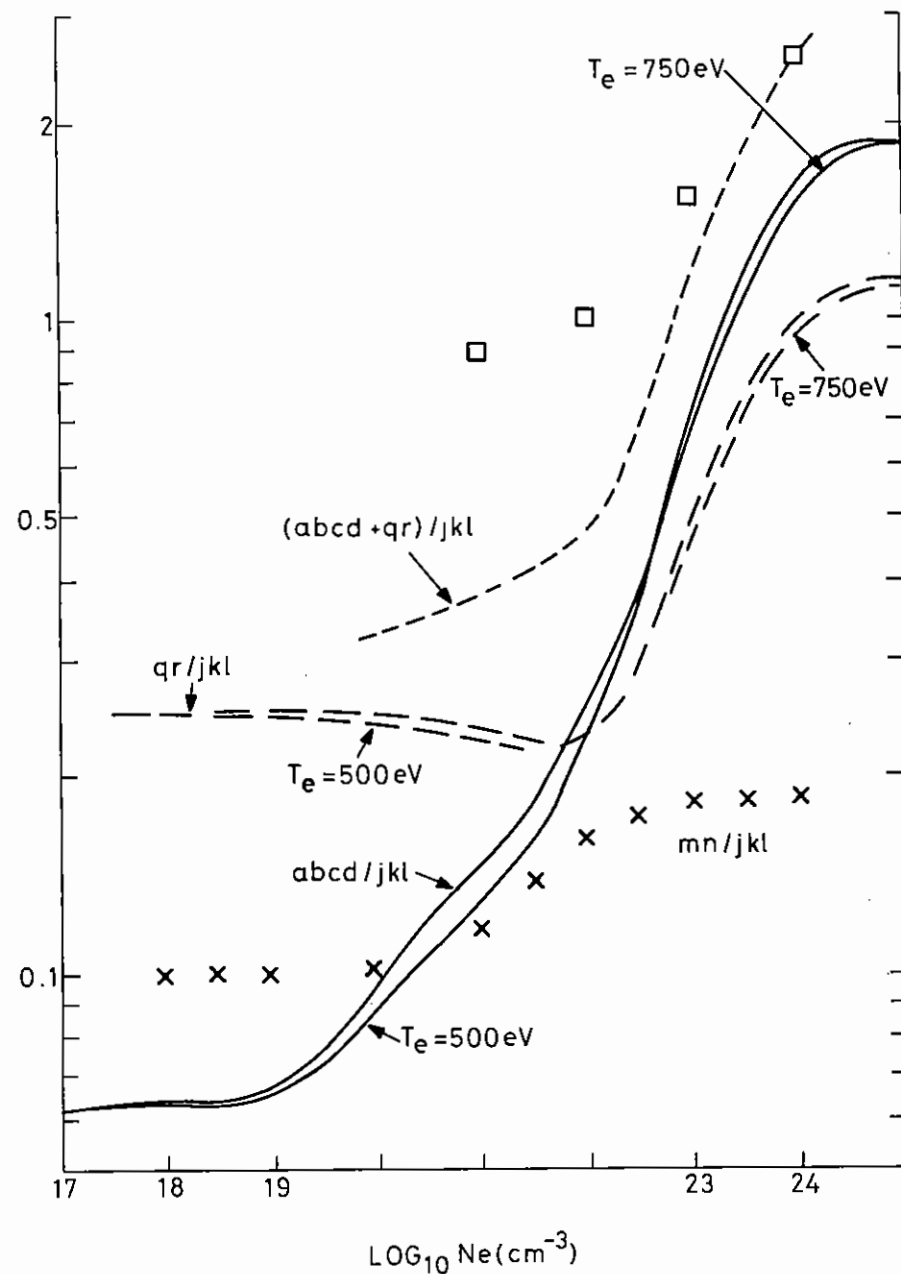


Fig. 7.15 The relative intensities of satellite complexes $abcd/jkl$, qr/jkl , mn/jkl and $(abcd+qr)/jkl$ as a function of electron density for Si. The variation with temperature is shown for $T_e = 500 \text{ eV}$ and $T_e = 750 \text{ eV}$. Results from Jacobs and Blaha are shown as $\square$.

It might appear from Figure 7.15 that the most suitable ratio for a density diagnostic is $abcd/jkl$. For Si, Figure 7.15 shows that this ratio is sensitive from 10^{20} to 10^{24} electrons/cm³. However below $n_e = 10^{22}$ cm⁻³ the variation of the ratio with n_e is weak, a factor of 1.5 for a factor of 10 variation in n_e . Moreover the spectral proximity of $abcd$ to the more intense qr line makes the $(abcd+qr)/jkl$ ratio easier to measure. Figure 7.15 shows there is a strong variation in this ratio only for $n_e > 10^{22}$ cm⁻³, when the ratio $R \propto n_e^{0.4}$ up to $n_e = 10^{24}$ cm⁻³.

The temperature variation of $(abcd+qr)/jkl$ is weaker. At about 10^{23} cm⁻³ the ratio scales as $T_e^{-0.1}$. Although inner shell excitation for Lithium-like Si is included (7.05), the variation of the ratio is still negligible.

The ratio of $(abcd+qr)/jkl$ taken from Jacobs and Blaha (7.05) is also shown in Figure 7.15. The discrepancy shown has caused the model in this paper to be scrutinised. At low n_e the results of this calculation agree with coronal limits (7.07) however the calculation contained in (7.05) does not appear to be converging to the coronal limit. Further there appears to be an error in the graphs of (7.05), as these values disagree with their tables.

Figure 7.16 shows the diagnostically simple ratio $(abcd+qr)/jkl$ for Na, Si, Cl and Ar. The density sensitive region increases with atomic number approximately as Z^6 , the inflexion point on Figure 7.16 occurring at an electron density n_e^* given by $n_e^* = (\frac{Z}{14})^6 \times 10^{23}$ cm⁻³. The scaling of the maximum ratio with density varies from 0.4 to 0.6 in Figure 7.16. The useful density range is approximately a factor of 20 for Na decreasing to about 10 for Ar. Calculations have not been carried out for any element beyond Ar because the jkl transitions split at higher Z and overlap the $abcd$ complex.

The ratio of the intercombination line ($1^1S - 2^3P$) to the resonance transition is principally a density diagnostic. However the problems are

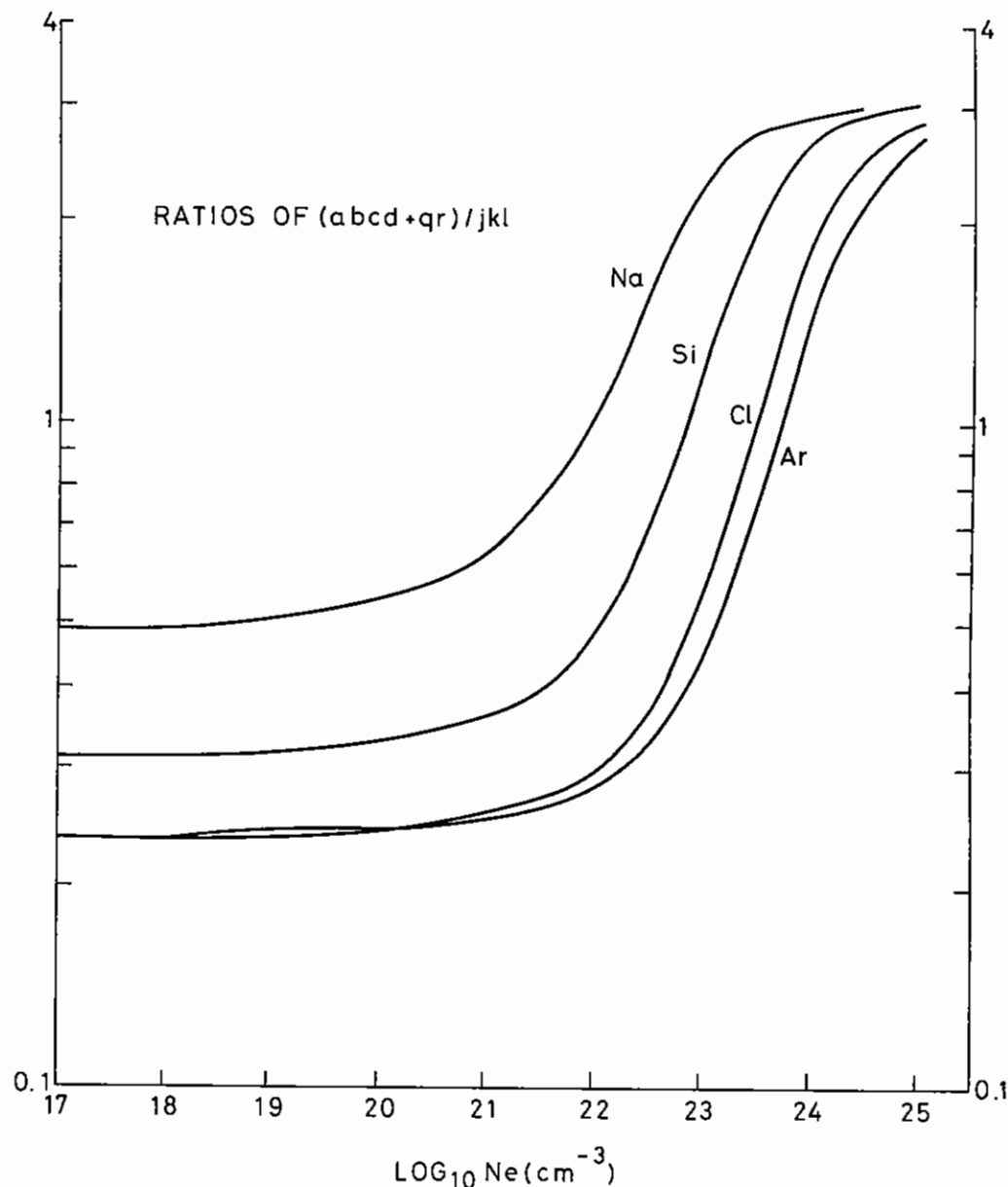


Fig. 7.16 The intensity ratio $(abcd+qr)/jkl$ for Na, Si, Cl and Ar at temperatures of 500 eV, 500 eV, 500 eV and 750 eV respectively.

firstly that there are underlying satellites mn and st that are difficult to resolve from the intercombination line and secondly that the $1^1S - 2^1P$ line is optically thick. The intensity of the satellites relative to the intercombination line can be evaluated from the code. To avoid opacity problems in measuring the $1^1S - 2^1P$ transition a higher member of the series can be used, although this introduces a slight temperature dependence to the ratio. The ratio of the $1^1S - 4^1P$ line to the intercombination line is shown in Figure 7.17 for the usual case when satellites are omitted, and for the case when the underlying st and mn satellites are included. As can be seen the satellites only make a significant change to predicted ratio at high density.

7.3.2(d) Experimental Results

(i) Ablation Plasma

Results were taken with 0.5 - 1.0 ns laser pulses producing ablation plasmas by irradiating NaCl and glass targets, with 0.53 μm radiation at an intensity of $5 \times 10^{14} \text{ W cm}^{-2}$. The foil targets were viewed with a time-resolving X-ray crystal spectrometer (7.26) using a flat Rubidium acid phthalate (RbAP) crystal and both a flat and curved Pentaerythritol (PET) crystal for Na and Si respectively. The time resolution was 100 ps. Data were recorded on calibrated photographic film and converted to film density and hence relative intensity by scanning a Joyce Loebel microdensitometer. A spectrum for Na is shown in Figure 7.18. It is notable that the resonance and intercombination transitions $1s^2 \ ^1S - 1s2p \ ^1P$ and $1s^2 \ ^1S - 1s2p \ ^3P$ are long-lived, ie 4 ns compared with the dielectronic satellites which are only emitted during the laser pulse. This clearly illustrates:

- (a) The advantage of the satellite lines as a density diagnostic as they record the highest density, and are only emitted during the laser pulse.
- (b) the inaccuracy of time integrated measurements especially of the resonance line.

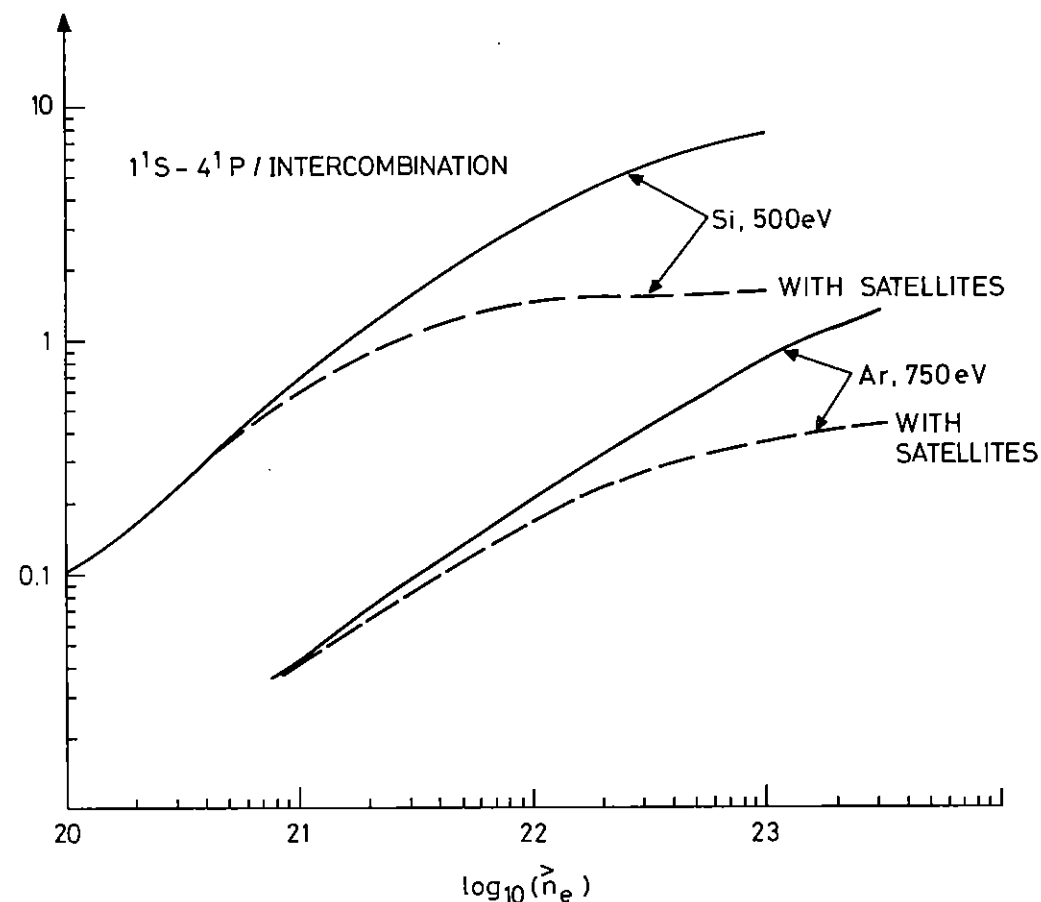


Fig. 7.17 The 1^1S-4^1P to intercombination line intensity ratio with and without the mn and st satellites included in the intensity of the 'intercombination' line.

Microdensitometer tracings of the spectrum are shown in Figure 7.18, 150 ps before and 150 ps after the peak of the laser pulse. (The line identifications for Figures 7.18, 7.20, 7.21, 7.22, 7.23 and 7.24 are given in Table 7.04.) At the earlier time the ratio of intensities

$$\frac{I(\text{abcd+qr})}{I_{\text{jkl}}} = 0.8 \pm 0.1$$

Figure 7.16 shows that the corresponding electron density is $0.4 \pm 0.1 \times 10^{22} \text{ cm}^{-3}$. The abcd complex is not clearly resolved from qr in Figure 7.18, but abcd appears to be slightly less intense than qr. Figure 7.19 shows numerical predictions for Na, and it can be seen that qr and abcd would have equal intensity at $0.3 \times 10^{22} \text{ cm}^{-3}$, which is consistent with their combined ratio to jkl above.

The tracing at late time in Figure 7.18 shows that the abcd+qr complex is substantially broader than jkl, giving

$$\frac{I(\text{abcd+qr})}{I_{\text{jkl}}} = 1.5 \pm 0.3$$

and from Figure 7.16 this indicates $n_e = 3 \pm 1.5 \times 10^{22} \text{ cm}^{-3}$. Correspondingly the tracing shows that abcd is more intense by possibly a factor of 1.6, which from Figure 7.19 indicates $n_e > 10^{22} \text{ cm}^{-3}$. The increase in the intensity of the satellites, with respect to the intercombination lines also indicates an increase in density at late time. Further the predictions are in concurrence with recently obtained density measurements made by time resolved Stark broadening which show densities of 10^{22} cm^{-3} .

Figure 7.20 shows the Si spectrum of an ablation plasma from a glass microballoon irradiated at $10^{14} \text{ W cm}^{-2}$ for 1 ns with $1.05 \mu\text{m}$ laser light. The use of a PET crystal in the spectrometer gave a greater dispersion than for the spectrum of Figure 7.18, and hence greater spectral resolution. The tracing and the print show the usual n222 dielectronic

TABLE 7.04

The identification of the lines on
Figures 7.18, 7.20, 7.21, 7.22, 7.23, 7.24

Identification on Figures	Transition
i	$1s^2 \ ^1S_0 - 1s2p \ ^1P_1$
ii	$1s^2 \ ^1S_0 - 1s2p \ ^3P_1$ and the mn and st dielectronic satellites
iii	qr dielectronic satellites
iv	abcd dielectronic satellites
v	jkl dielectronic satellites

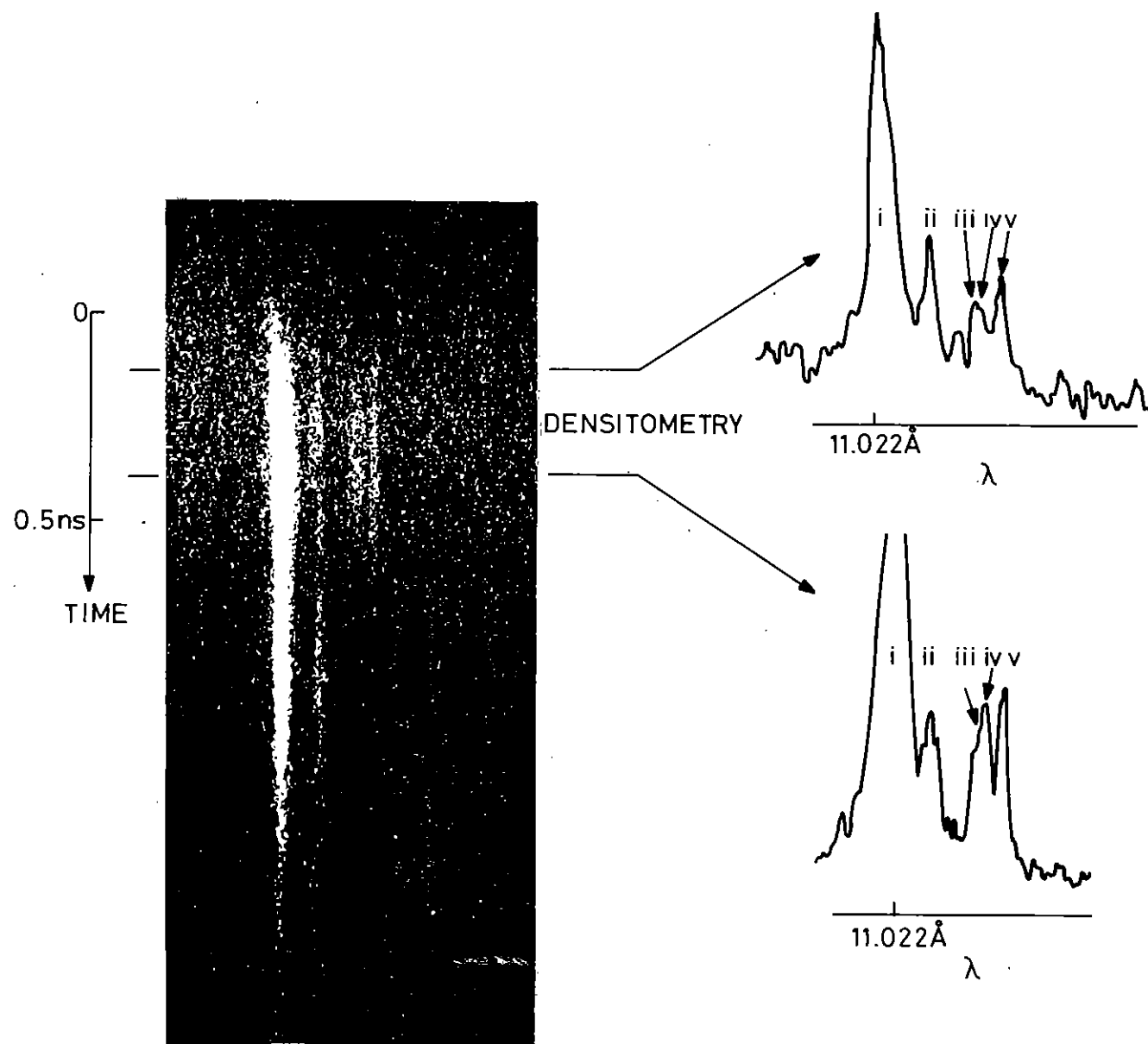


Fig. 7.18 A time-resolved spectrum of the He-like 1 - 2 complex in Na from a NaCl plasma. The laser pulse lasts for the length of the continuum emission. Microdensitometer tracings at early and late times in the laser pulse are shown. The identifications are given in Table 7.02.

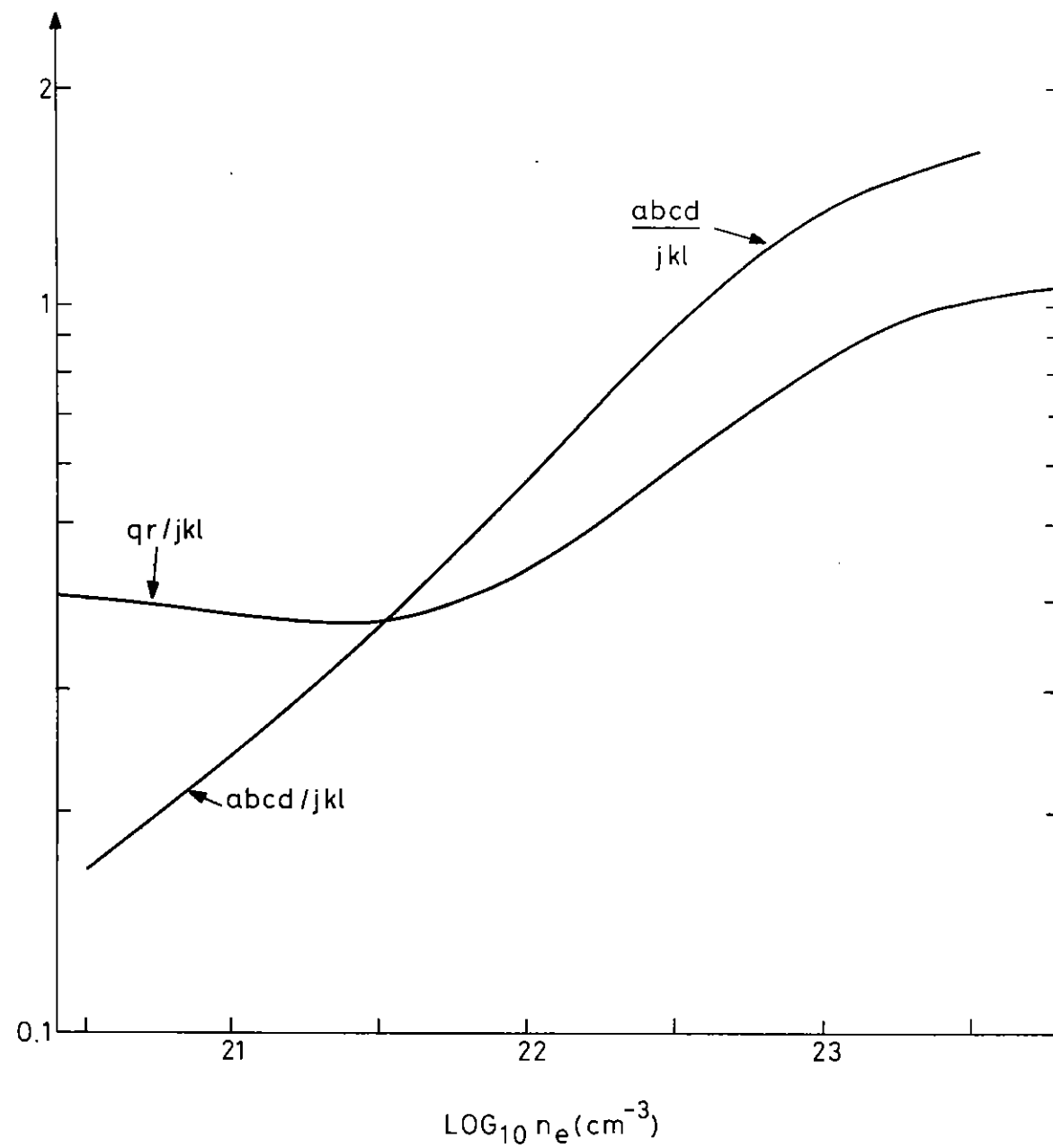


Fig. 7.19 Model predictions of the intensity of the ratios of the $\frac{abcd}{jkl}$ and the $\frac{qr}{jkl}$ complexes in Na as a function of electron density at 500 eV.

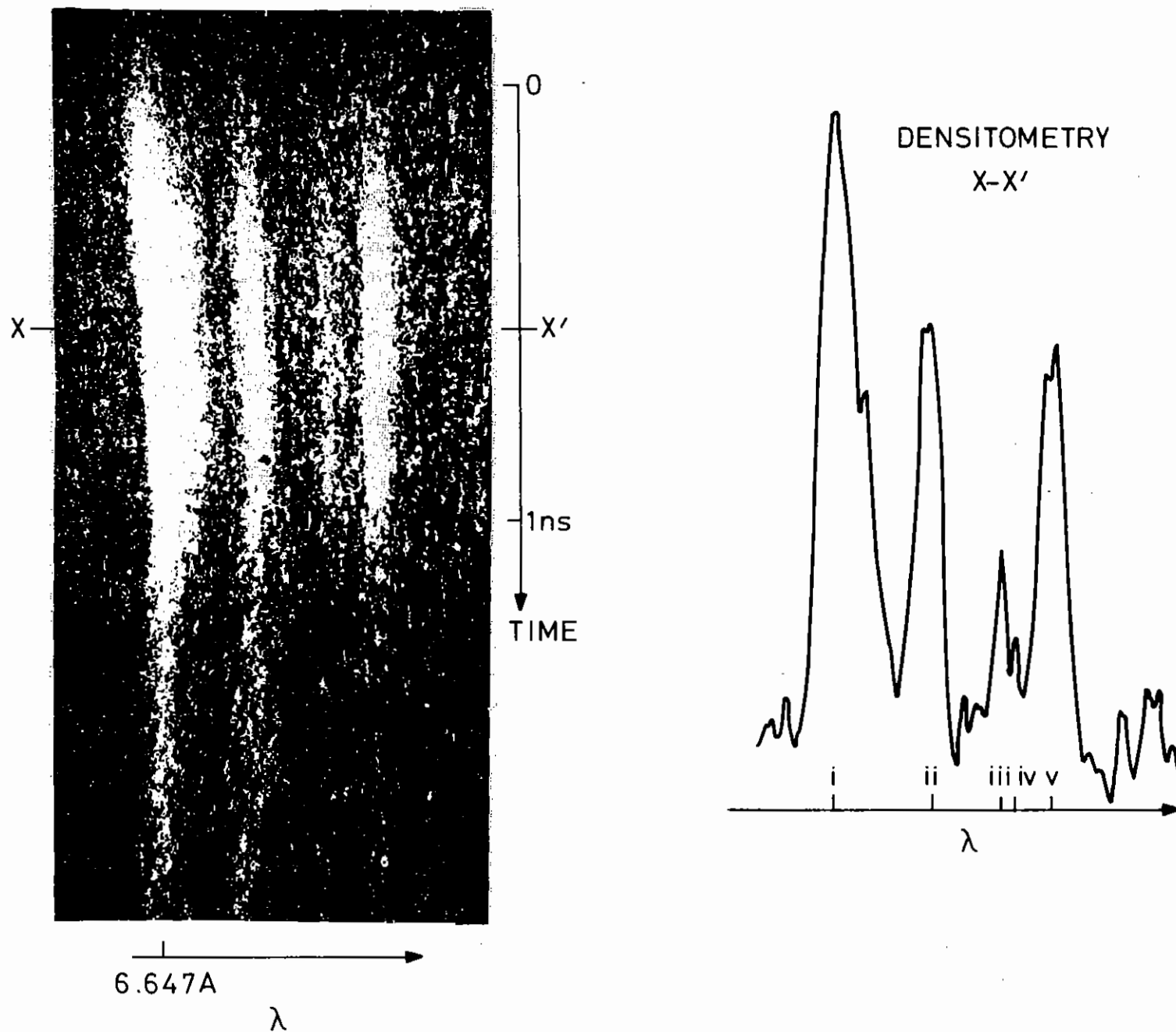


Fig. 7.20 Time-resolved spectrum of Si ablation plasma with 1.05 μm irradiation of a glass microballoon. The laser pulse is 1 nsec and because of the longer wavelength has a lower critical density than in Fig. 7.18. Consequently the satellite ratio is not a useful density diagnostic (see text).

satellites and unresolved satellites close to the resonance line, presumably due to $n = 3$ spectator electrons (7.28)(7.13).

The $n = 2$ satellites give a lower electron density than for the $0.53 \mu\text{m}$ irradiation of NaCl , Figure 7.18, as would be expected for the lower critical density, viz 10^{21} cm^{-3} for $1.05 \mu\text{m}$ irradiation, cf $4 \times 10^{21} \text{ cm}^{-3}$ for $0.53 \mu\text{m}$ irradiation.

The densitometer tracing shows

$$\frac{I(\text{abcd+qr})}{I(\text{jk}\ell)} = 0.30 \pm 0.1$$

which, from Figure 7.16 indicates an electron density $n_e < 2 \times 10^{21} \text{ cm}^{-3}$. This is too low for the ratio to be an effective diagnostic in Si. The abcd complex is not well resolved but gives

$$\frac{I(\text{abcd})}{I(\text{jk}\ell)} \approx 0.16$$

Figure 7.15 would then give $n_e = 2 \times 10^{21} \text{ cm}^{-3}$.

Ablation plasmas were also produced by irradiating plane silicon targets with $0.53 \mu\text{m}$ laser light at $5 \times 10^{14} \text{ W cm}^{-2}$.

Silicon spectra are shown in Figure 7.21 with densitometer tracings at two different times. It is notable that at the later time the ratios are higher, that is

$$\begin{array}{ll} \frac{I(\text{qr})}{I(\text{jk}\ell)} = 0.42 & \frac{I(\text{abcd})}{I(\text{jk}\ell)} = 0.2 \\ \text{cf earlier } \frac{I(\text{qr})}{I(\text{jk}\ell)} = 0.16 & \frac{I(\text{abcd})}{I(\text{jk}\ell)} = 0.065 \end{array}$$

Figure 7.15 then shows that for silicon this indicates $n_e < 10^{21}$ at early time, increasing to $2 \times 10^{22} \text{ cm}^{-3}$ at later time. However inspection of Figure 7.15 shows that despite the higher density than in Figure 7.20 Si

is only just becoming viable as a density diagnostic. Figure 7.21 illustrates that although in principle the ratio $\text{abcd}/\text{jk}\ell$ is a density diagnostic for very low n_e , in practice the proximity of abcd to qr, which is constant and larger than abcd at low density, makes it difficult to use as a low level density diagnostic.

(11) Compressed Core Plasmas

Higher electron densities have been achieved with so-called exploding pusher plasmas. In these targets densities can be varied from about $1 \times 10^{23} \text{ cm}^{-3}$ up to $5 \times 10^{23} \text{ cm}^{-3}$ for thick shell targets or $0.53 \mu\text{m}$ implosions. A time-integrated spectrum from an exploding pusher plastic target filled with Freon 12 (CCl_2F_2) and consequently emitting X-ray lines from Helium-like Cl, is shown in Figure 7.22. The abcd complex is not well resolved from qr in this case and

$$\frac{I(\text{abcd+qr})}{I(\text{jk}\ell)} = 1 \pm 0.1$$

which from Figure 7.16 shows that $n_e = 3.0 \pm 0.5 \times 10^{23} \text{ cm}^{-3}$. The electron density can also be obtained from the Stark broadening of the Helium-like Cl $1s^2 \ ^1S - 1s4p \ ^1P$ line on this shot. Comparison with theoretical calculations (7.29) gives $n_e = 1.5 \pm 0.5 \times 10^{23} \text{ cm}^{-3}$.

In contrast, the emission from a low density Cl plasma is shown in Figure 7.23. This plasma, as in Figure 7.18, was produced by irradiating plane NaCl targets by $0.53 \mu\text{m}$ laser light at $5 \times 10^{14} \text{ W cm}^{-2}$. In Figure 7.23 the helium-like chlorine lines are shown. The ratio

$$\frac{I(\text{abcd+qr})}{I(\text{jk}\ell)} = 0.2$$

which, from Figure 7.16, implies that the density is not high enough for the ratio to be useful as a diagnostic, ie $n_e < 10^{22} \text{ cm}^{-3}$.

Finally, Figure 7.24 shows the spectrum from a high density argon-filled exploding pusher irradiated at $0.53 \mu\text{m}$. The abcd+qr complex is 5% more

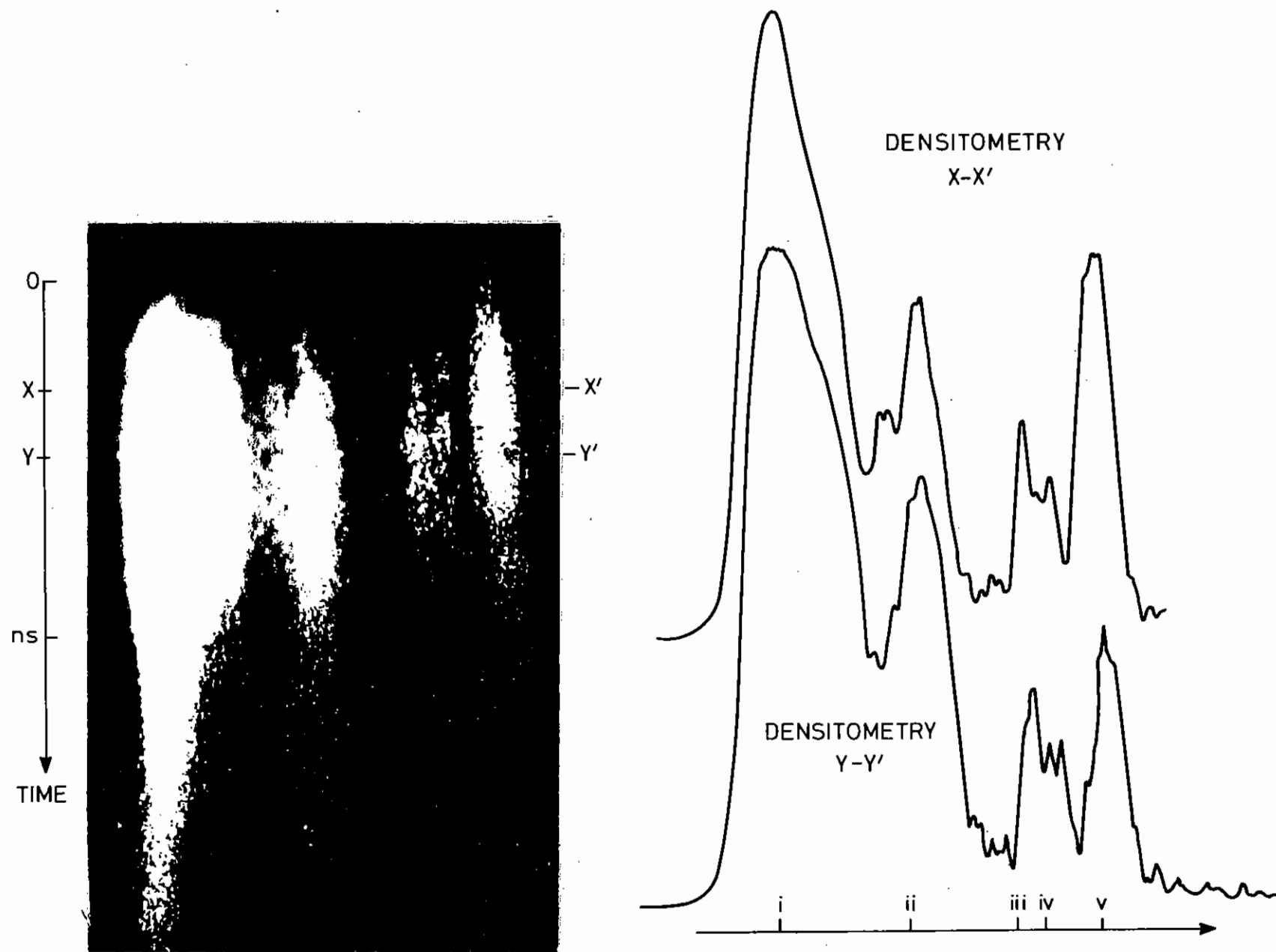


Fig. 7.21 Time-resolved spectrum of a Si ablation plasma produced by irradiating a planar SiO_2 target with an $0.53 \mu\text{m}$ laser beam at $5 \times 10^{14} \text{ W cm}^{-2}$. The spectrum indicates a higher density than Fig. 7.20.

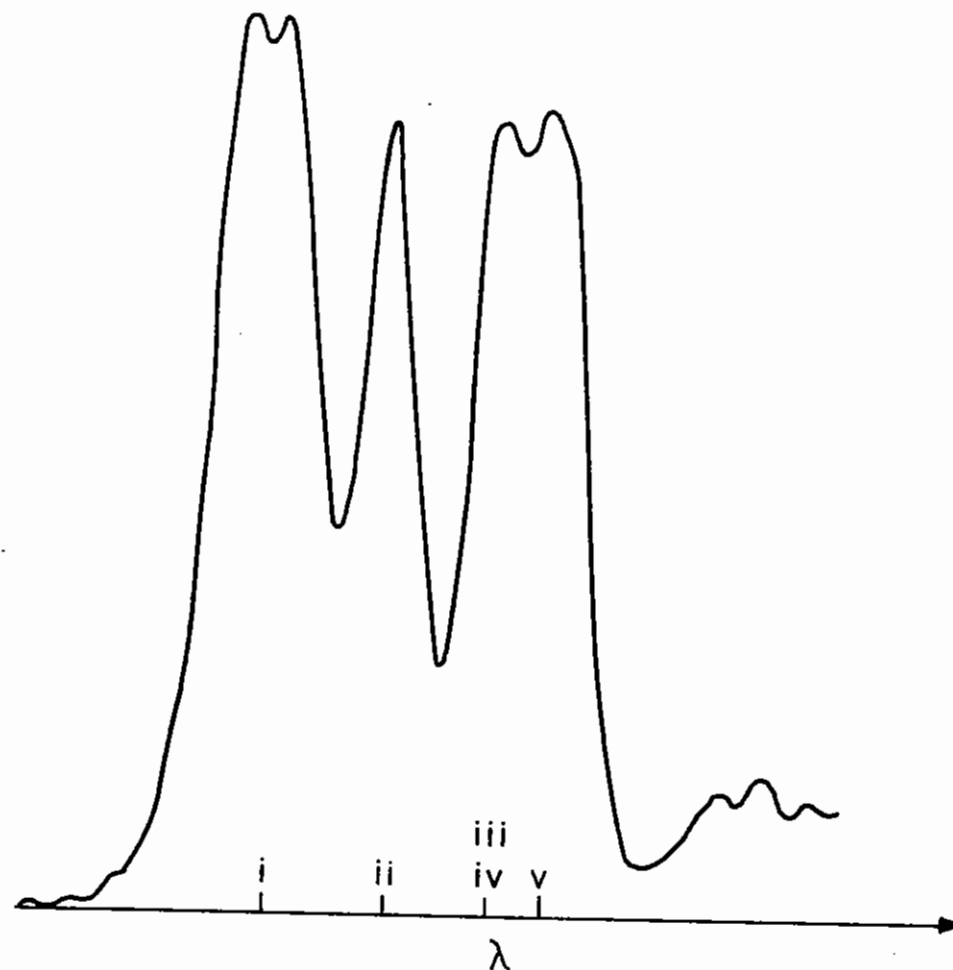


Fig. 7.22 Time-integrated spectrum showing the helium-like 1 - 2 complex in CII from an exploding pusher plasma. This was produced by symmetrically irradiating a plastic microballoon filled with 1 Bar of Freon 12, with 80 J of 1.05 μm laser light in a 100 psec pulse.

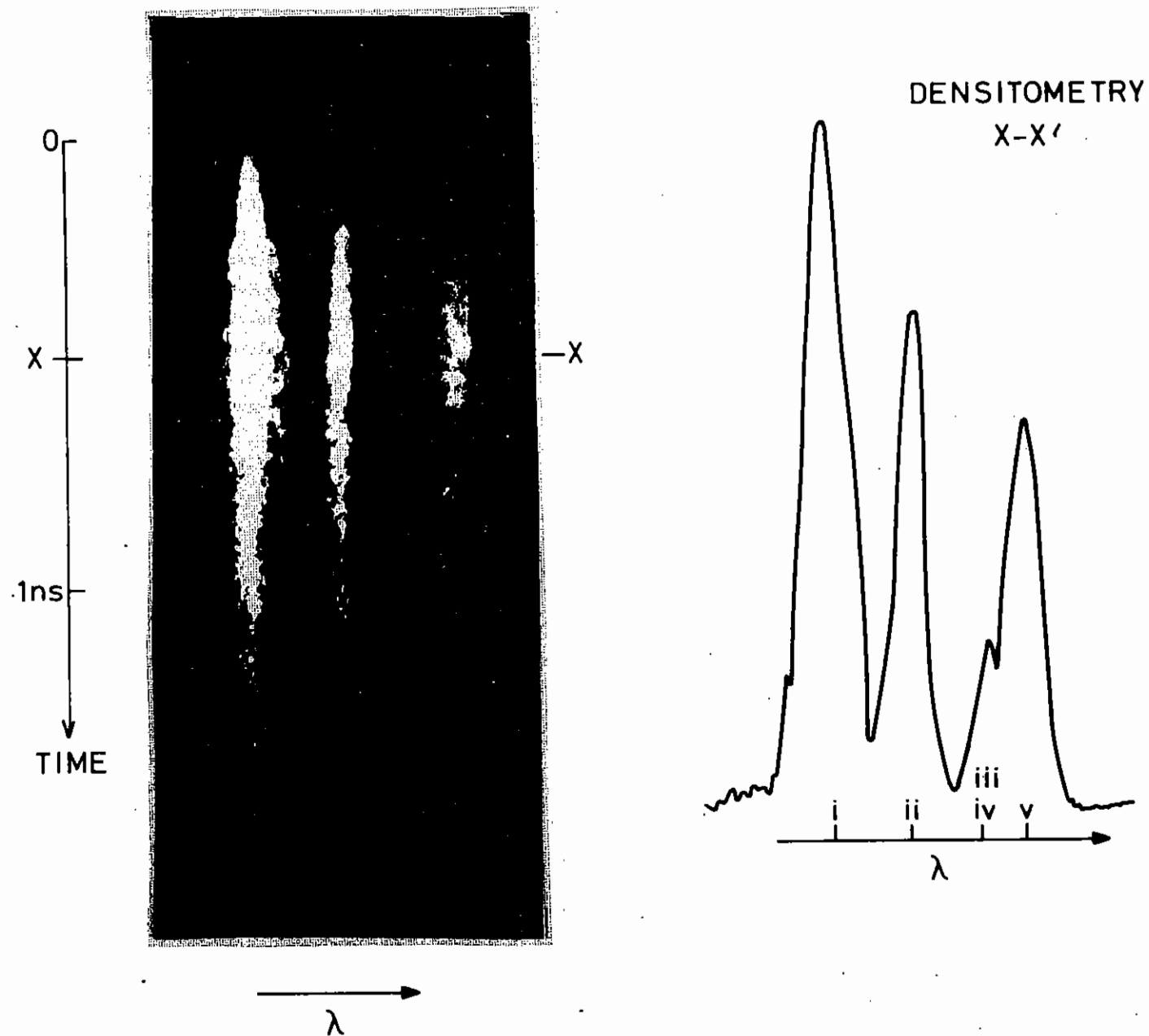


Fig. 7.23 Time-resolved spectrum showing the helium-like 1 - 2 complex in Cl from an ablation plasma. This was produced by irradiating a NaCl planar target with a single $0.53 \mu\text{m}$ laser beam at an intensity of $5.0 \times 10^{14} \text{ W cm}^{-2}$.

intense than $jk\ell$, indicating $n_e = 5 \times 10^{23} \text{ cm}^{-3}$. It is also notable that the dielectronic satellites are more intense than the intercombination line at these densities. Line broadening on this shot indicates $n_e = 4 \times 10^{23} \text{ cm}^{-3}$.

7.3.2(e) Discussion and Conclusion

The use of the He satellite ratios as a density diagnostic depends on accurate prediction of the collisional cross-sections between the autoionising levels and the Li-like species and, less importantly, on the ionisation balance. Direct experimental verification of the collision rates is difficult; however there is little reason to doubt the accuracy of the theoretical predictions. The ionisation balance has also to be either known, or calculated. In the laser plasmas described here, long laser pulses (ie 1 ns) have been used, allowing sufficient time for a collisional radiative equilibrium to be set up in the ablation plasma.

It has been pointed out by some workers (7.05)(7.30) that inner shell ionisation is necessary for accurate predictions of the satellites to Lyman-like transitions. These are implicit in the present calculation.

The present calculation shows that the satellite ratio $abcd+qr$ to $jk\ell$ is essentially a density diagnostic with a useful dynamic range of 10 for $Z = 10$ up to 20 for $Z = 18$.

Predictions of the code are verified by results from six different experiments. Experiments on ablation plasmas with $n_e = 10^{22} \text{ cm}^{-3}$ show that Na can be used as a density diagnostic, whereas the atomic number of Si is too high for it to be particularly useful and that of Cl is much too high. The Si data indicate some consistency in that the densities indicated by the $abcd/jk\ell$ ratio are similar to the densities indicated by the $qr/jk\ell$ ratios.

Results from higher density exploding pusher targets, where Stark broadening has been used to corroborate the density independently, are also presented. for densities $2 \times 10^{23} \text{ cm}^{-3}$, Cl is used as a density



Fig. 7.24 Time-integrated spectrum of the He-like 1 - 2 complex in Ar. This was produced by irradiating argon-filled glass micro-balloons with $0.53 \mu\text{m}$ laser light.

diagnostic, and at $4 \times 10^{23} \text{ cm}^{-3}$, Ar is used.

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7.4 K-Edge Fiducials to see Absorption Shift in Resonance Line Positions

A preliminary experiment has been carried out in order to absolutely measure the spectral position of resonance lines emitted from laser compressed plasmas. Argon filled exploding pusher targets produced 500 eV - 800 eV cores of density about 10^{23} electrons per cc. The density was measured by fitting the experimental line profiles to the theoretical line shapes generated from the optically-thin Helium-like profiles of R W Lee with LTE opacity broadening (7.31). Measurement of associated line shifts due to the high density was attempted using the K absorption edge of Potassium as a spectrometer calibration. Figure 7.25 shows an Argon core spectrum taken through a $1.5 \mu\text{m}$ thick layer of KCl. The proximity of the Potassium edge to the Argon XVII $1s^2 - 1s3p$ (He β) resonance line coupled with a knowledge of the spectrometer dispersion allowed determination of the line position to an accuracy of $\pm 0.7 \text{ m\AA}$.

Figure 7.26 shows the digitized experimental He β line compared with the theoretical line position at an electron density of $1.6 \times 10^{23} \text{ cm}^{-3}$.

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7.5 Quantitative X-ray Measurements for Radiative Preheat Determination

The absolute X-ray emission of laser irradiated SiO_2 microsphere targets has been measured in the 0.9 to 1.5 keV region. Recent theoretical (7.32) and experimental study (7.33) of $\lambda = 0.53 \mu\text{m}$ ablatively driven microshell implosions has pointed to X-radiation from the laser interaction region as a possible source of target preheat. Preheating of the shell during acceleration will cause it to decompress resulting in reduced dynamic

pressure and a poorer implosion.

Some work on measuring radiation production has been published (7.34) but the bulk of the data is for $1.06 \mu\text{m}$ irradiation of planar targets and none at all is directly applicable to $0.53 \mu\text{m}$ irradiation of glass microspheres.

The measurements performed at Rutherford Appleton Laboratory were obtained using 20 - 50 J of green light in three beams to illuminate a hemisphere of a microballoon target. The three beams were incident along three mutually orthogonal axes with the spectrometer collecting X-rays from the point on the sphere midway between the three beams. Only three of the usual six green beams were available at the time of these measurements. The f/1 beams were defocused by 3.5 target radii beyond the target centre producing an average irradiance on the region of the target observed by the spectrometer equal to:

$$I_{\text{inc}} \approx \frac{E_{\text{inc}}}{3\pi r^2} \frac{1}{\tau_L} \text{ W cm}^{-2}$$

τ_L , the laser pulse width at full width half maximum power was measured as 900 ps.

The spectrometer used was a Mark 4 minispectrometer described in previous reports (7.35). A calibrated TLAP Bragg diffracting crystal ($2d = 25.75 \text{ \AA}$) was used to produce a spectrum on No Screen X-ray film which has been extensively calibrated (7.36). A $5.5 \mu\text{m}$ thick Aluminium filter was placed over the spectrometer slit in order to shield the crystal from debris from the target. A $5.5 \mu\text{m}$ filter was used despite its low transmittance (6% at 925 eV) in order to strongly attenuate Silicon resonance lines which were observed in second order. The weak second order lines served as a useful wavelength calibration but were not bright enough to affect the first order continuum as they are quite close to the K absorption edge of Aluminium and were thus severely attenuated.

The observed spectral intensity I_e is shown in Figure 7.27 for the five

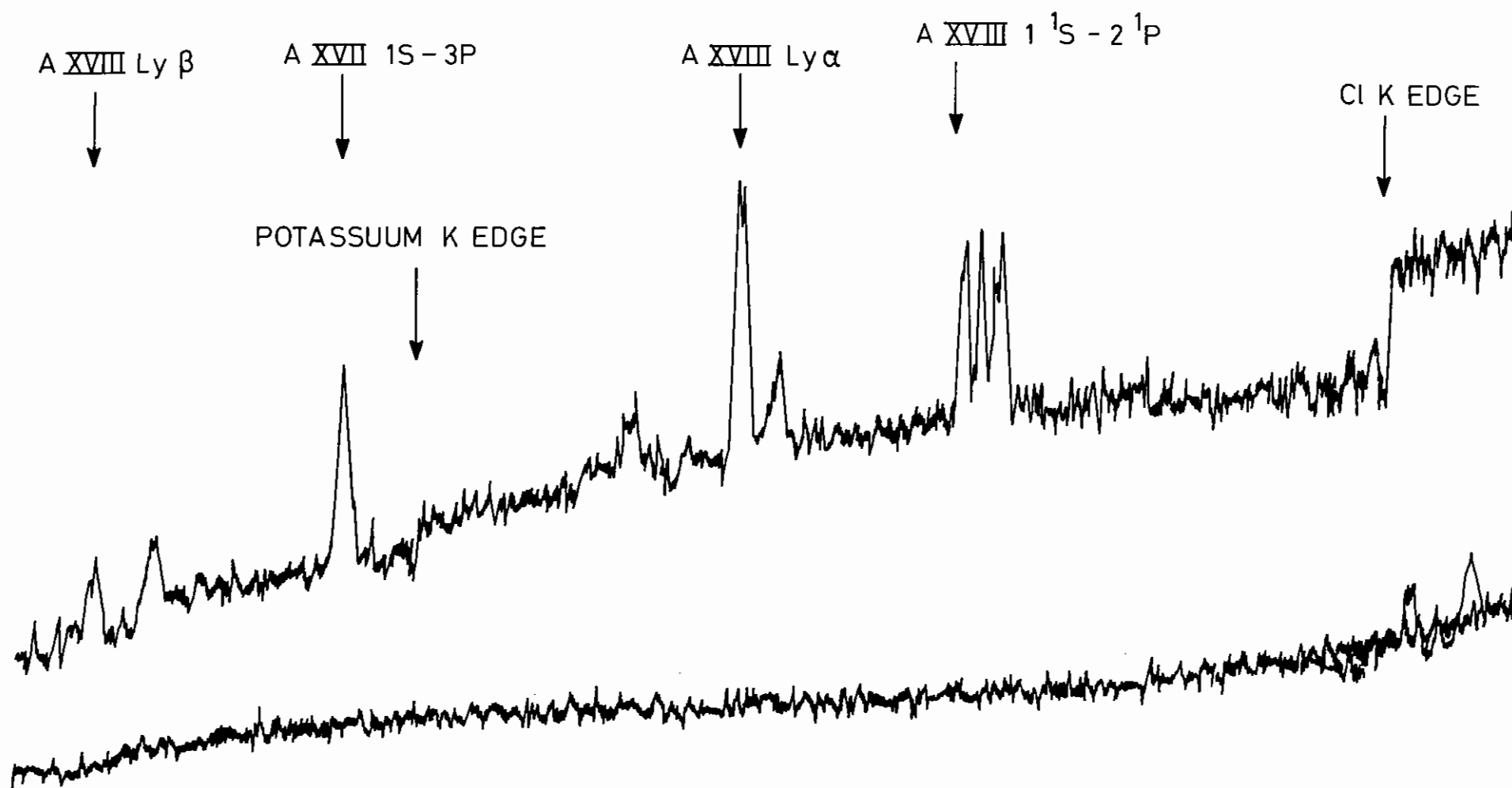


FIG.7.25. ARGON EMISSION SPECTRUM THROUGH A COLD K Cl ABSORBER TO
FIX THE ABSOLUTE WAVELENGTH SHOT 11160682

targets irradiated. Details of the targets are shown in Table 7.05. Simulations (7.32) indicate a peak in the spectrum in the region of 500 eV. Extrapolations of the spectra to 500 eV indicate total conversion fractions of 4 - 12% of incident energy into X-rays for SiO₂ targets, depending on the spectral form assumed for I_e below 900 eV. For plastic shell targets (CH) the conversion fractions were down by a factor of 2.5 and were in the range 1.5 - 5%.

The main value of these measurements was to enable normalisation of the radiation transport package (7.32) included in the one dimensional MEDUSA simulation. The code was in quite good agreement with the measurements and reproduced the sharp fall-off in I_e ($\sim \exp\{-h\nu/140 \text{ eV}\}$) in the region 900 eV - 1.5 keV.

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7.6 Time dependent Stark Broadening measurements in Planar-Laser Target Experiments

7.6.1 Introduction

The interest in plasmas produced by laser irradiation of a planar target has increased due to the relation of these plasmas to direct-drive laser compression studies; to a fundamental interest in the ablation plasma characteristics; and finally to the desire to perfect diagnostic techniques for transient high density plasmas. For these planar target plasmas are convenient for their ease of observability and simplicity.

The methods to predict density from these laser plasmas have relied on spectroscopic information such as line ratios, line widths, and scattering measurements, (7.37) (7.38). The scattering measurements are limited to densities much less than critical density, i.e. $< 4.10^{21}$ for a frequency doubled Nd laser, and therefore do not address the problem of the physical conditions in the critical density region. The scattering measurements, however, have been time and space resolved. The spectroscopic methods, which rely on spectrally resolved data have tended

TABLE 7.05

SHOT	TARGET MATERIAL	DIAMETER (μm)	E _{inc} (J)	I _{inc} (Wcm^{-2})
09 191182	SiO ₂	175	51	$7.8 \cdot 10^{13}$
24 231182	SiO ₂	110	24	$9.3 \cdot 10^{13}$
08 231182	SiO ₂	106	21	$8.8 \cdot 10^{13}$
16 181182	CH	136	37	$9.5 \cdot 10^{13}$
12 181182	CH	176	29	$4.4 \cdot 10^{13}$

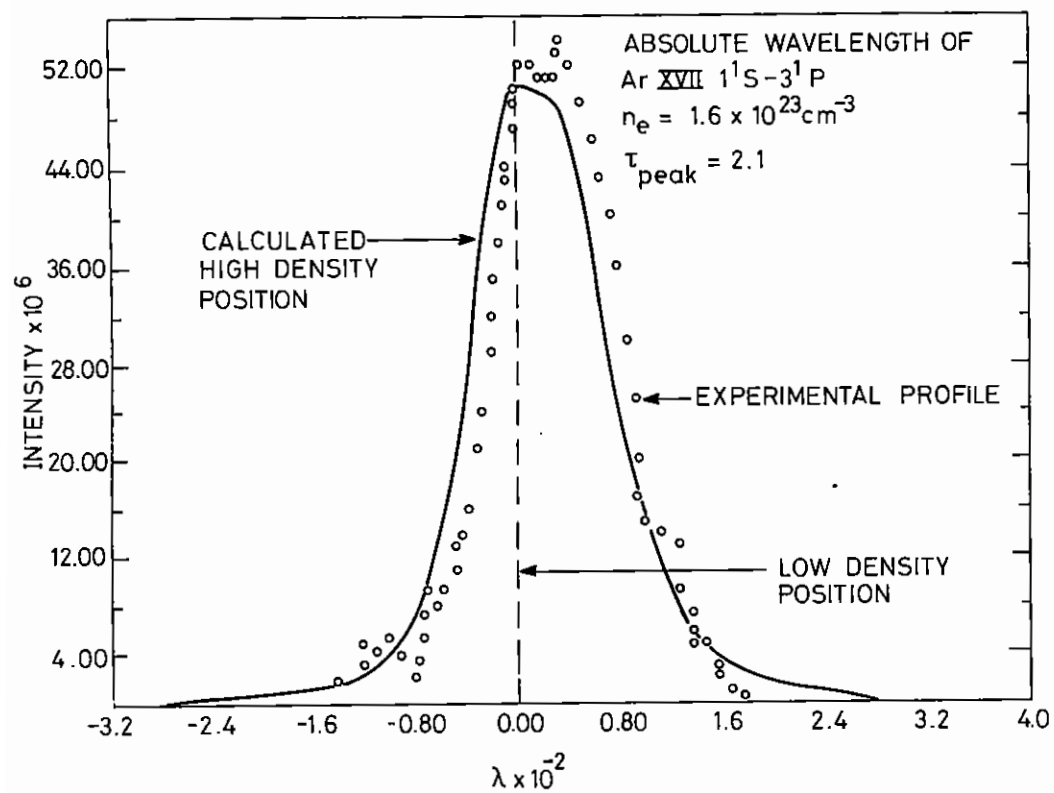


Fig. 7.26 Digitised experimental He β line compared with the theoretical line position at an electron density of $1.6 \times 10^{23} \text{ cm}^{-3}$.

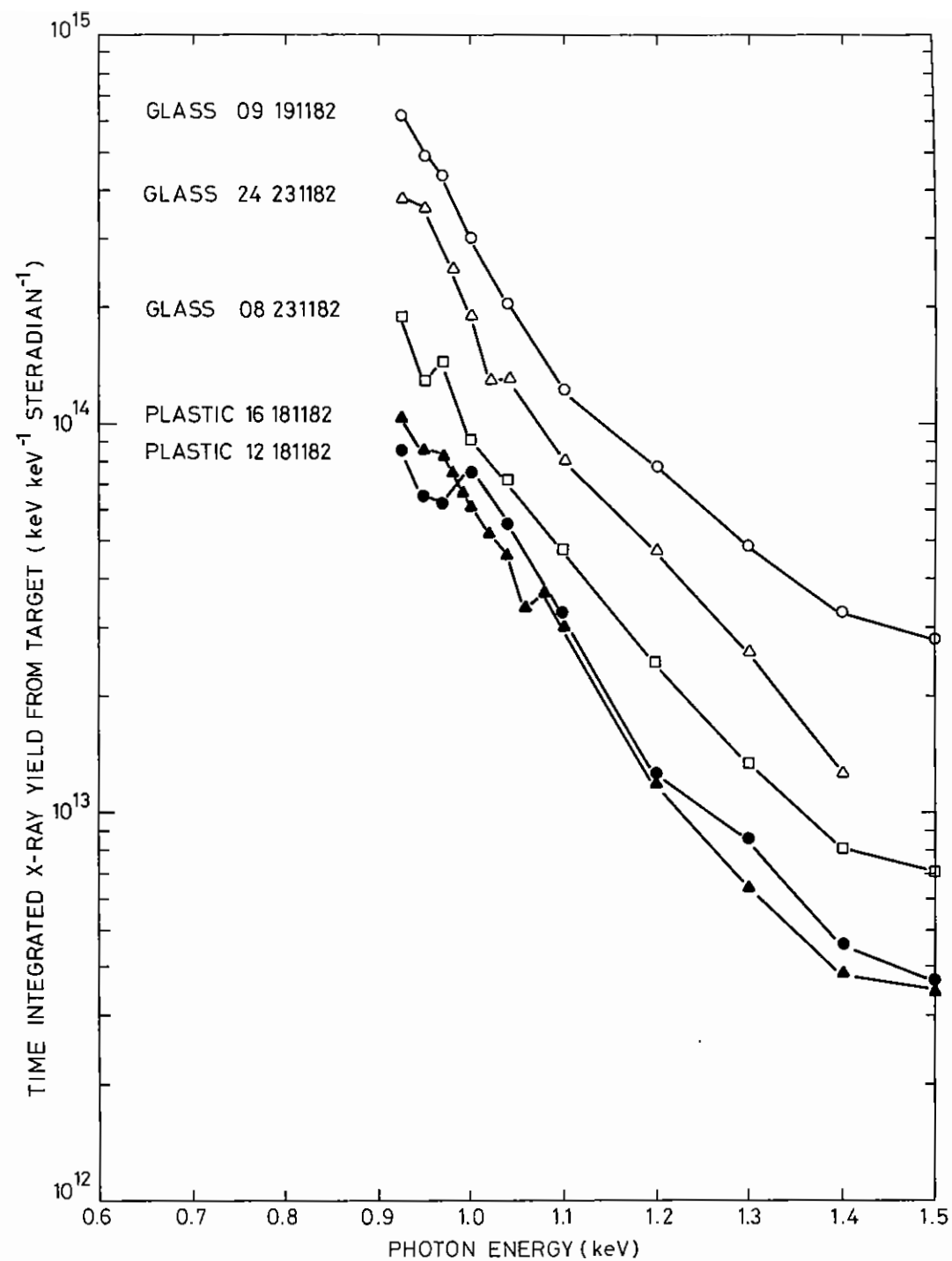


Fig. 7.27 Absolute x-ray production from the three beam irradiation of microsphere targets. The estimated systematic error on each point is $\pm 30\%$.

to be time integrated, but space resolution has been obtained (7.39). The basic argument to justify time integrated data is that a time averaged result represents the system.

The difficulty with these spatial integrations can be surmounted by using a small amount of the element of interest in a localized spot (7.39). This tracer will then be localized in the plasma and can give some spatial information about a localized position in the plasma. However, this still requires assumptions concerning the time dependence of the plasma to extract information. Thus one major step to improve spectroscopic diagnostic procedures is to use time resolved spectroscopy to determine the temporal variations of the electron density in the region about the critical density. In the present work, we report on the use of an x-ray crystal streak spectrograph to study the time-dependent electron density of an ablation plasma by using the Stark broadening of the higher principal quantum number states from the Lyman series of hydrogenic ions.

This work goes some way to completing the necessary diagnostics studies which will eventually lead to a time resolved, space resolved experimental knowledge of these ablation plasmas. The present experiment which has time resolution of 80 psec does not have spatial information. However, the present work is conclusive proof that streak camera data can be used in a spectrally resolving mode to produce line profile information.

7.6.2 Experimental design

The experiments were performed at the Central Laser Facility of Rutherford Laboratory in a collaborative effort between Imperial College, London, the Laser Program at Lawrence Livermore National Laboratory, and the Rutherford Laboratory. The experiments used the single beam target chamber where one beam of frequency-doubled $1.05\mu\text{m}$ laser was focused by a f/l lens onto planar targets. The pulse length was nominally 900 psec and the average intensity was $\sim 10^{15}\text{W/cm}^2$ (7.40).

The experimental diagnostics, beside the standard laser beam diagnostics, consisted of an x-ray streak camera with a crystal as a dispersing

element. The observations were made with the target tilted 10° toward the streak camera. There was also a time-integrated, space-resolving crystal spectrometer viewing the target parallel to its surface and a pinhole camera.

The choice of material used was dictated by many factors. As a starting point, we know from prior experiments that the canonical temperature would be $T_c \sim 1\text{keV}$ and the peak electron density would be approximately the critical density, which for the frequency doubled Nd laser is $N_e \sim 4 \cdot 10^{21}\text{cm}^{-3}$. The various factors which lead from the T_e , and N_e to the choice of Al can be delineated. First, one requires a reasonable population in the hydrogenic levels of the species chosen. The population can be estimated theoretically using a steady state rate equation model (even though the actual level populations depend on time) to assure this. Next, the members of the Lyman series must be sufficiently intense to be observed and yet remain optically thin. This can be estimated by reference to line profile calculations together with calculated level populations to yield the opacity and emissivity, while plasma dimensions can be primarily obtained from pinhole picture and/or space resolving spectra (see Table 7.06). The line profile must then be primarily broadened by the Stark effect so that the widths of the profiles are dominated by the plasma density. This implies that the Stark width must be greater than the total Doppler and instrument width. This again requires a line broadening calculation. Further, the plasma density will cause the merging of the high n-state members of the Lyman series and will depress the ionization potential. These effects will obscure the profiles of interest, thus one must ensure that the transitions are observable - simple estimates of these effects will suffice (see Table 7.07). Finally, the emission will arise from plasma with velocity differences relative to the observer. Therefore, we minimize this effect on the width of the profile by choosing the broadest lines possible. These considerations lead to a choice of element about $z \sim 13$.

For this experiment the targets were pure Al disks. A spectrum observed with a streak camera as an example of the emission from the blow-off

TABLE 7.06

Estimate of Lyman Series Opacity for $\lambda = 100 \mu\text{m}$, $T_e = 750 \text{ eV}$.

n_{upper}	$\phi (\text{cm}^2)$	τ
2	$.56 \cdot 10^{-16}$	112
3	$.29 \cdot 10^{-17}$	5.8
4	$.10 \cdot 10^{-17}$	2
5	$.21 \cdot 10^{-18}$	.42
6	$.13 \cdot 10^{-18}$	.26
7	$.40 \cdot 10^{-19}$	.08

TABLE 7.07

IP depression (IPD) and Inglis-Teller Limit (ITL); The Highest principal quantum number existing vs electron density.

n_{MAX}	IPD $N_e (\text{cm}^{-3})$	n_{MAX}	ITL $N_e (\text{cm}^{-3})$
16	10^{19}	27	10^{19}
11	10^{20}	18	10^{20}
8	10^{21}	12	10^{21}
7	$4 \cdot 10^{21}$	10	$4 \cdot 10^{21}$
7	$6 \cdot 10^{21}$	9	$6 \cdot 10^{21}$
6	10^{22}	9	10^{22}

plasma observed is shown in Fig. 7.28. Note that the spectral range covered, from about 2.1-2.4 keV, is restricted due to the cathode size and the high spectral dispersion. The lines here, the Al XIII ls-np transitions where $n = 5, 6, 7, 8$ are observed for a time duration of about 1600 ps. The laser pulselength is ~ 800 ps on this shot.

To analyze the results obtained from the experiments, the film is digitally scanned using a densitometer. Spectral scans in various time intervals are taken and corrected for film response and dispersion, producing a relative intensity vs x-ray energy spectrum. A similar analysis is performed on the space resolving spectrometer data.

In Fig. 7.29 an example of the analyzed streak camera data is shown for the time interval averaged between 1100 and 1200 ps. Note the Lyman transition $n = 5$ to 1 through $n = 7$ to 1 are clearly observed. Spectra of this type illustrate our basic results and are analyzed to yield the time history of the plasma.

The shots recorded were taken over a period of time and represent two different sets of conditions. In one set we have a short laser pulse length ~ 300 psec half width and an energy on target of about 20 joules. The second set has a longer pulse length of about 1000 psec with 30 joules on target. These distinctions are seen in the analysis of the electron density time histories as will be shown below.

7.6.3 Theoretical Modeling

To analyze the data we require a computational model that can produce the line shapes of the Lyman lines for high n states. This is straightforward in that the Lyman series lines can be calculated for all interesting n , given that a microfield has been calculated and that account is taken of the overlapping of the higher n lines. The overlapping is incorporated by using a schematic cutoff for the minimum momentum transfers allowable. This cutoff includes both the Debye wave number and the wave number necessary to give rise to overlapping lines. This approximation together with the fact that the width operator decreases logarithmically beyond the plasma frequency and that it should decay as $\Delta\omega^{-1/2}$ beyond the Weisskopf

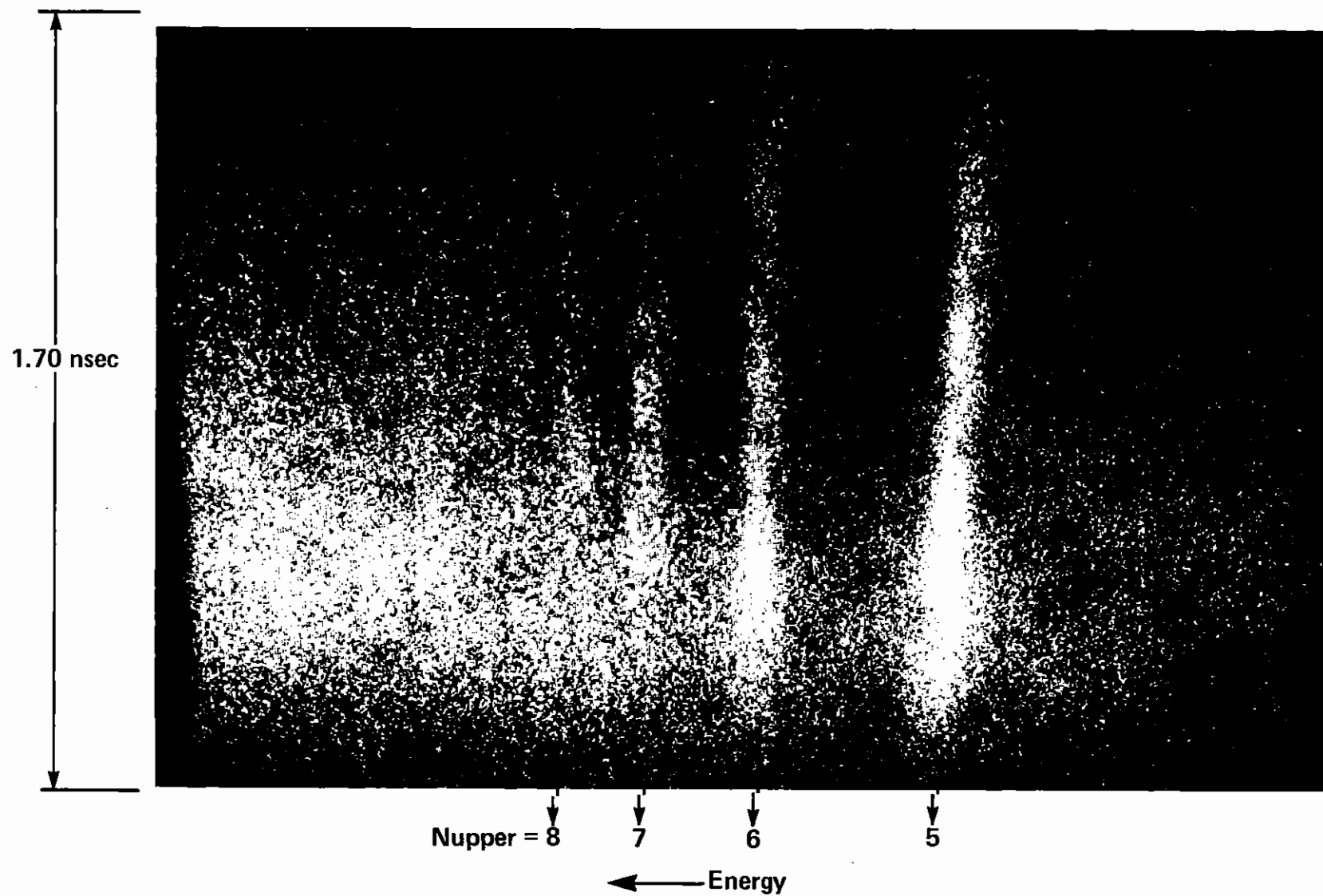


Fig. 7.28 The x-ray streak camera spectrum for shot CLF 1216S7. The time history of the x-ray emission of an Al XIII Lyman series. The lines are denoted by the upper principal quantum number and the time is in nanoseconds.

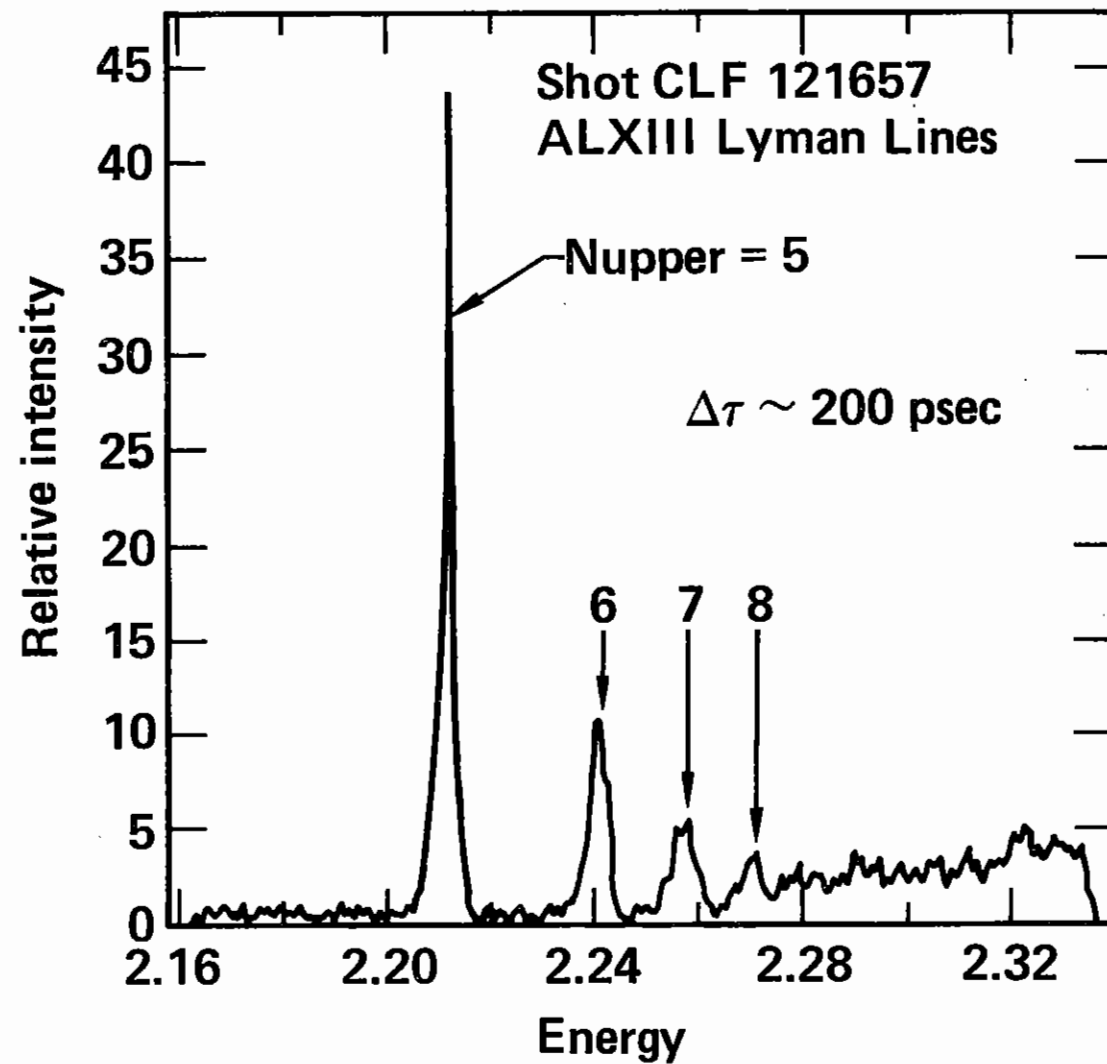


Fig. 7.29 A "line out". The spectrum averaged over the time interval 1.1 to 1.2 ns from shot CLF 1216S7. The intensity is relative and the energy scale is in keV. The Al XIII Lyman lines are noted.

frequency yields a shift operator which compares well with detailed calculations using a relaxation theory type model for the width operator (7.41) (7.42).

The total line profiles compare well with previous published line profiles, or widths.(7.41) (7.42). The comparison for hydrogenic ion lines gives less than a few percent difference over the entire range of electron densities previously published. Therefore, it is felt that the accuracy of these schematic calculations is comparable to those performed in greater detail.

The results of the calculations for the Lyman $n = 4, 5, 6, 7$ lines of Al XIII are shown in Fig. 7.30. Here we have plotted with full-width at half maximum in eV of the lines without Doppler effect versus electron density in cm^{-3} . The temperature for a planar laser-target plasma. However, since Doppler convolution is not included and since the widths are relatively insensitive to temperature this plot can be used in a wide variety of cases.

Also included in Fig. 7.30 are the widths of the $n = 4$ to 1 and $n = 6$ to 1 lines measured from the "shoulders". That is, these profiles have a central unshifted peak, which is quite narrow, and a broader shoulder region. It is useful to have these "shoulder widths" because the central components are prone to the highest uncertainty from ion-acoustic modes, ion dynamics, Doppler shifts, and opacity corrections. Further, experimental measurements of the central narrow peak width is also difficult since it requires high resolution. Thus, since the shoulder width of the transition 1-4(6) is essentially the same as the total width of 1-5(7) one may have twice the diagnostic information.

To analyze the experimental data, we require total profiles which include Stark plus Doppler and instrument broadening. The instrument broadening was derived from estimates of source size, crystal performance, and streak camera resolution, and this leads to a half width at half maximum of .75 eV for the energy range of the Al XIII lines studied here. The Doppler broadening is determined from the temperature diagnosed using time

resolved line ratios of Kauffman, et al (7.41). The nominal temperature of 1000 eV is found to be high and we will use 750 eV for our study.

In Fig. 7.31 we have plotted the full width at half maximum vs electron densities of the same spectral lines shown in Fig. 7.30. Here we have included all the broadening mechanisms. As can be seen in Fig. 7.31, the lines with central dips ($n = 4, 6$) behave as predicted in Fig. 7.30. However, the lines with central spike show a non-monotonic behaviour due to the fact that the Doppler and instrument half-half width $\Delta\omega_{ID}$ is constant (by assumption) as the electron density varies. More explicitly, if we take the $n=4$ to 1 transition as an example we see three distinct regimes which occur as density increases.

$$\begin{array}{ll} \bullet \Delta\omega_{ID} > \Delta\omega_S & n_e < 10^{20} \\ \bullet \Delta\omega_S > \Delta\omega_{ID} > \Delta\omega_\epsilon & 10^{20} < n_e < 5 \cdot 10^{21} \\ \bullet \Delta\omega_\epsilon > \Delta\omega_{ID} & n_e > 3 \cdot 10^{23} \end{array}$$

where $\Delta\omega_S$ is the shoulder width, and $\Delta\omega_\epsilon$ is the width of the central spike. Thus we have a good diagnostic from all lines $n = 4$ through 7 in the electron density region from 10^{20} cm^{-3} and to about 10^{22} cm^{-3} .

The result provides us with the necessary information to analyze the experimental data. Curves similar to those shown in Fig. 7.31 are the primary diagnostic tool. Further description of the theory are contained in other works for example, see Lee, et al., 1982, Lee and Whitten, 1983.(7.44) (7.45). It is important to note that the accuracy of this method is very high since, for example, our accuracy in measurement is better than .25 eV in half width and as a worst case this yields about 25% inaccuracy in the critical density region.

7.6.4 Results

Comparison of theoretical and experimental line shapes can be carried out as follows. We fit the individual lines to determine, from the width, the

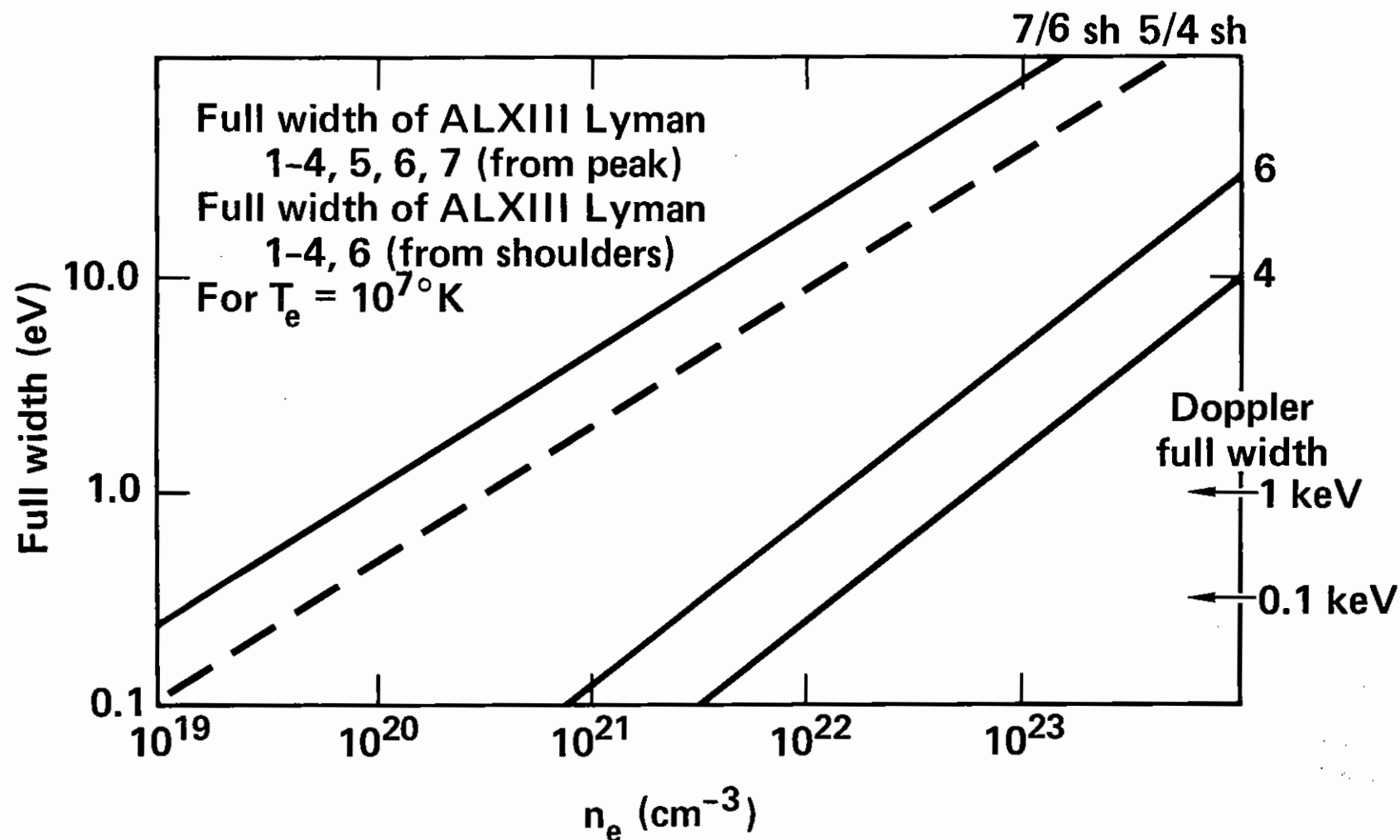


Fig. 7.30 The theoretical full width at half maximum in eV versus electron density in cm^{-3} . The results are shown for Al XIII Lyman 1-4, 5, 6, 7 and for the "shoulder widths" of 1-4, 6 (see text). The temperature is fixed at 1 keV for the results shown. The Doppler widths corresponding to 1 keV and 0.1 keV are also shown.

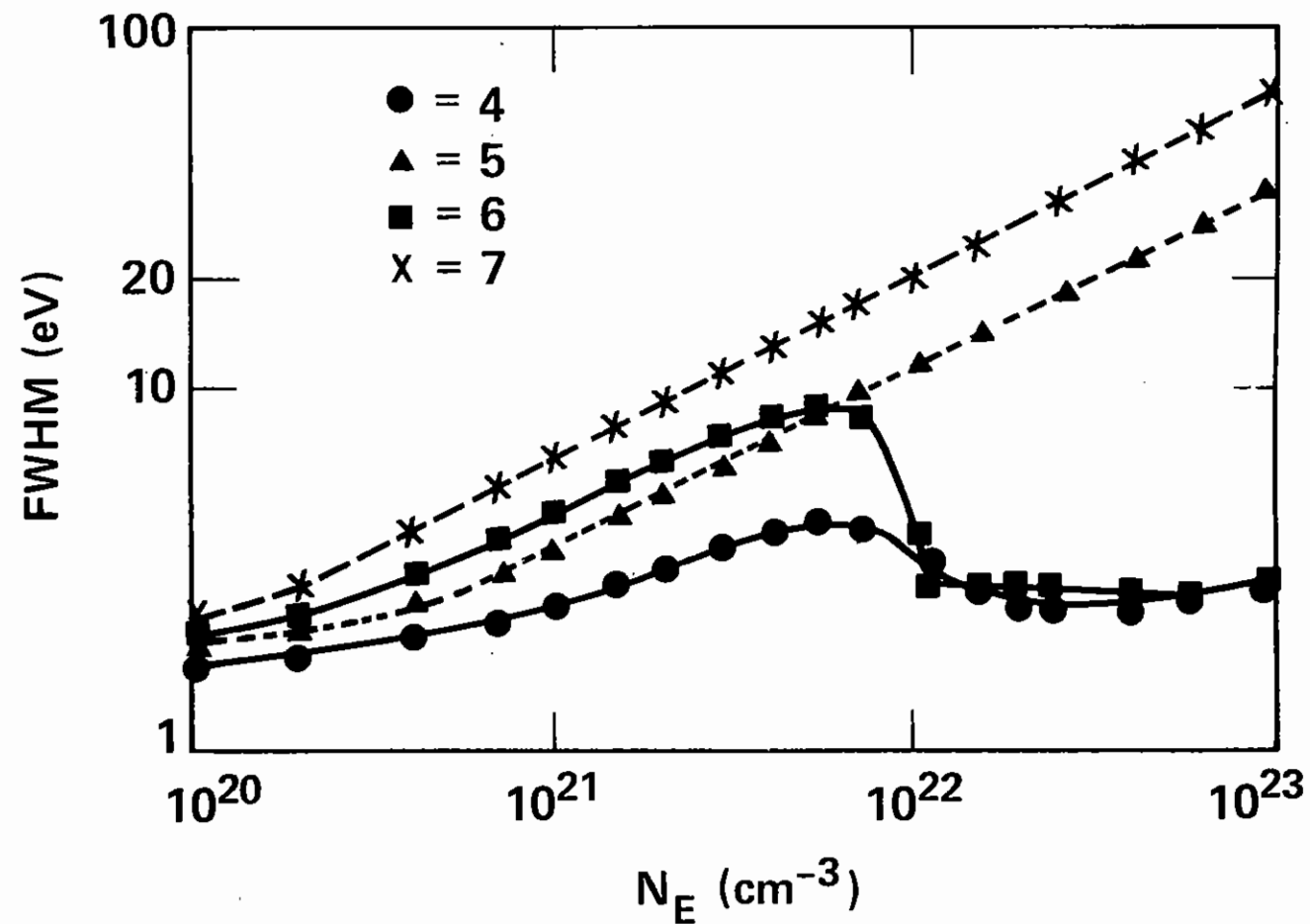


Fig. 7.31 The FWHM vs electron density for the Lyman transition of Al XIII with all the broadening mechanisms. Doppler width at $T_e = 750$ eV. Instrument width at .75 eV (HWHM).

electron density. Since we have data for at least two lines -- and sometimes four lines -- this information should produce a good indication of the electron density as a function of time.

As indicated earlier, we have data from two sets of experiments. To illustrate the differences we show in Figs. 7.32 and 7.33 examples of the electron density vs time for the two sets of data. In Fig. 7.32 we have results of CLF1217S7, which shows a fast rising n_e to a peak of about 600 psec after the image appears on the film. (Note that the times here are relative and are only compared to the time when the first emission is seen). The maximum electron density is $\sim 2 \cdot 10^{21} \text{ cm}^{-3}$ and the half-width of the electron density spike is about 400 psec. This behaviour is corroborated by the fact that we do have information from three different transitions. Further there is a plateau region after 900 where the electron density stays fairly constant. To better understand this behaviour we indicate in Fig. 7.32 the time history of the intensity of the continuum. This gives an indication of the laser pulse characteristics. It can be seen that the density follows the intensity and the plateau in density is produced by the laser heating of the expanded under dense plasma. This indicates that for a period the plasma emission arises from a region of approximately critical density and for an approximately equal period, the emission arises from a region in this case which is ~ 1.4 critical density. This indicates that attributing one density from a time independent width is insufficient. Further the time averaged electron density which is determined by time integrating the streaked data are shown in Fig. 7.32 by arrows. Note that a simple average of the electron density over the time of observation produces average $0.95 \times 10^{21} \text{ cm}^{-3}$ while the electron density predicted from the time integrated profiles is about 30% higher. This seems to be good agreement, but also shows that the weighting is towards higher density as one expects (see 7.46).

The other set of data taken is represented by Fig. 7.33. This shot, CLF1201S14, shows some interesting differences when compared to Fig. 7.32. First note that a streak camera was set to cover the $n = 4$ to 1 transition and higher lying lines. This provides us with further information in this

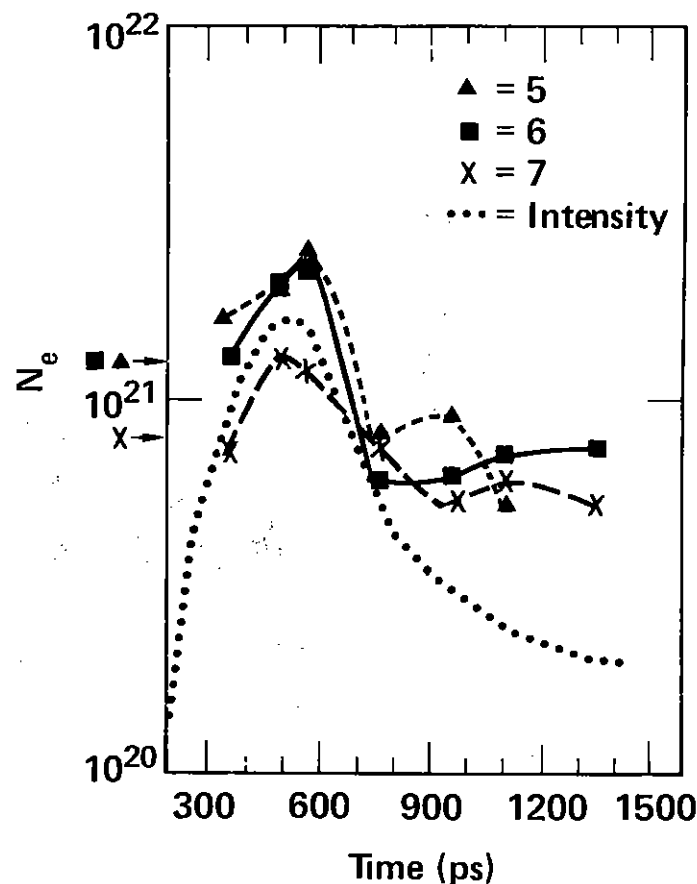
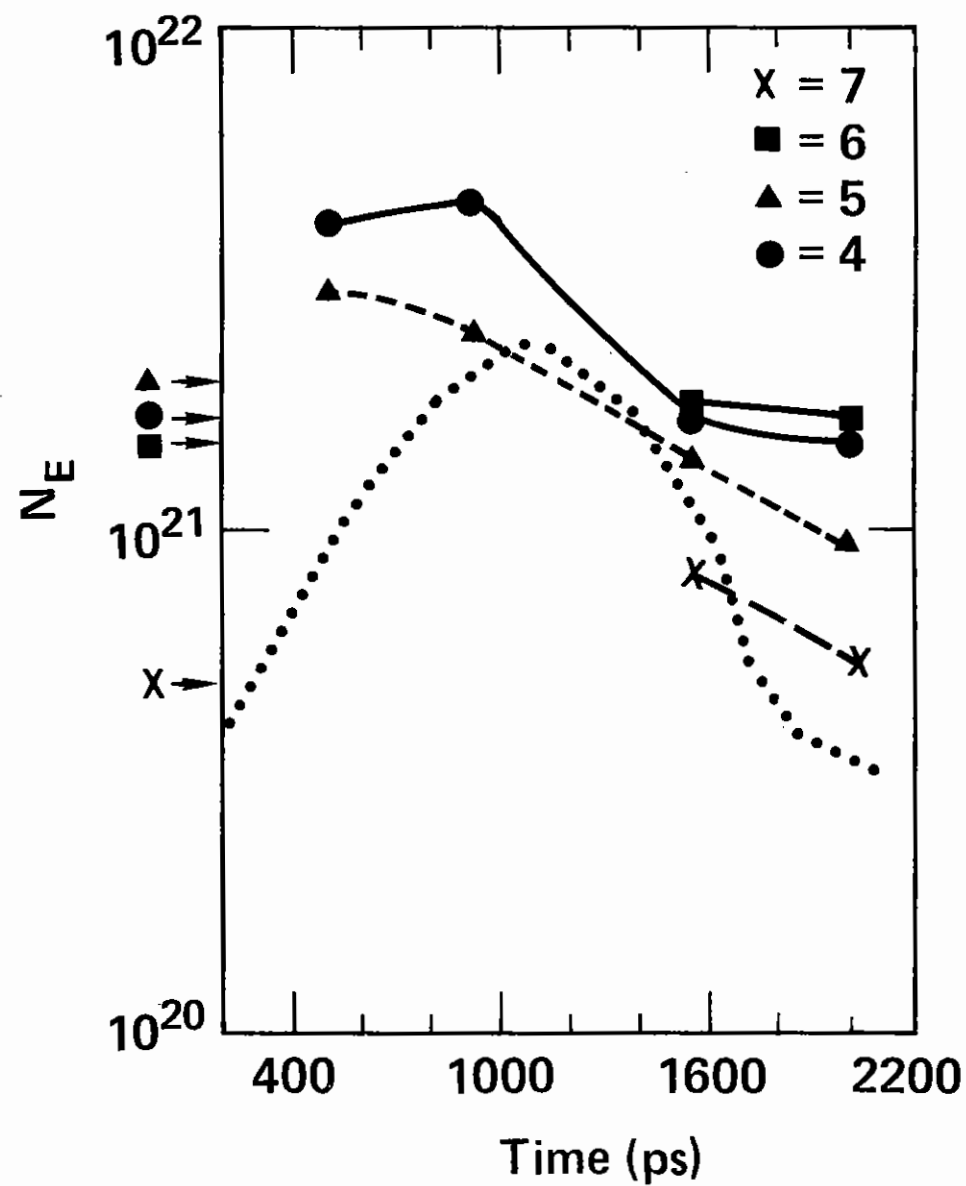


Fig. 7.32 The predicted electron density in cm^{-3} vs time (ps) for shot CLF 1217S7 for three transitions.
 o - - - - o = transition 5 to 1
 o - - - - o = transition 6 to 1
 x x = transition 7 to 1
 The arrows indicate the time integrated electron density.
 I I = relative intensity.



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Fig. 7.33 The predicted electron density in cm^{-3} vs time (ps) for shot CLF 1201S14 for four transitions as Fig. 7.32 plus $\square\square\square$ = transition 4 to 1

experiment. Next, the peak electron density is higher in this case and the general trend is to have smoother variation of the electron density with time. The intensity also shows a smoother variation. Further, the mean electron density as determined from the time integrated profiles is clearly weighted toward the lower electron density. This is in distinction to the average of the electron density which is generally higher than that predicted by the time integrated profiles. Note that the average density tends towards the value 1.10^{21} which is an interesting point given that this value seems to represent much of the data taken previously in time integrated experiments.

The major differences between the two sets of data therefore are that the peak electron density reached is different and the time history follows the different pulses. The occurrence of a different peak electron density is not yet fully understood. A possible interpretation is that the longer pulse length experiments which have a longer rise time can create a longer scale high density region in front of the critical density surface. This high density, combined with critical to produce emission which is interpreted as arising from "critical density". In the shorter pulse length experiments the high density region does not spread away from the critical density surface. This leads to emission which is produced only at critical density and lower. Therefore, we interpret the mean density as arising from less than critical.

7.6.5 Discussion

There are a number of comments which can be made concerning the data and analysis procedure employed to obtain the electron density. First there are a number of issues which concern the theory for Stark broadening of these transitions.

In the theory the effect of overlapping lines, of different principal quantum numbers, are treated as separate entities. In the case of these high n lines at high density, the effect of these Stark crossings should be calculated in detail. Next, the central portion of the profiles with unshifted Stark components, i.e., the transitions $n = 4$ and 6 to 1 , are

the most liable to theoretical error. The difficulties with the central components have been discussed previously by Griem et al (7.43), and Lee (7.42).

However, the fact that these profiles may require further theoretical development does not seem to affect the validity of the model in this experiment. Thus, although the need for detailed studies of the effects of ion dynamics, ion acoustic modes, and radiative transfer is obvious, these do not adversely affect the interpretation of the data in these experiments. It should also be noted that the line spectra observed potentially provides enough information to determine temperature from a study of the relative intensities of the line series member and the continuum. While this is possible in principle, it requires a more complete theoretical model which is under study presently.

The major problems with interpretation of the data concerns other subsidiary effects which can be estimated and are found to be either unimportant in the present study or give indications of new experimental procedures which correct the situation. The physical effects are ionization potential depression, line merging, and opacity, while the experimental effects are the laser intensity effects (intensity, pulse length and spot size) and effects of gradients.

The estimate of the opacity can be made assuming the levels reaches a steady state - which at the highest densities found here is a valid assumption - and taking a worst case condition. We assume that the number in the ground state of the hydrogenic species is half of the total ion density, n_I , which is assumed to be $(Z-1)n_e$, and the electron temperature is 750eV . For these parameters the peak opacity per number density per length, ϕ , of each line in the Lyman series of Al XIII is given in column 2 of table 7.06. Thus the opacity is determined by $\tau = \phi(n_I * .5) * \ell$, where ℓ is the geometrical length of the system. In the present case we estimate $\ell = 100 \mu\text{m}$ to give an upper bound on τ . As seen in Table 7.06, column 3, these upper bound (7.41) on the opacities indicate that $n_{\text{upper}} > 4$ should always be optically thin, while care should be taken with the $n = 4$ to 1 transition. For the present analysis, therefore, opacity is not a factor.

The problem with observability is compounded by both IP depression and line merging. The effect of both these processes is to cause the high n lines to disappear into the continuum. Since we are observing the high n states this is indeed a problem. In Table 7.07, the highest principal quantum number which should exist (even though it would be close or at the continuum) is given for various electron densities. By reference to the table the more critical effect is seen to be the IP depression. To arrive at the IP depression, we have used a simple ion sphere model to estimate the depression which for these densities it is appropriate. The point is that for the experimental set which is represented by the data in Fig 7.41, we have indirect evidence that the transitions above $n = 6$ are obscured into the continuum. This seems to be an indication that the IP depression and line merging criteria underestimate the value of the upper level principal quantum which is observable. That this is an underestimate is corroborated by a more complete model which includes line broadening, merging, and IP depression consistently (see Lee et al, 1982, and Lee and Whitten, 1983 (7.44) (7.45)). These physical processes are seen to effect the data and therefore must be considered in any experimental design. However, they do not modify the analysis done here.

The final comments are reserved for the experimental conditions. It is important to note that the density follows the laser intensity history. Therefore to achieve a better understanding of these ablation plasmas we need to ensure that laser parameters are controlled. In this way we can test diagnostic performance in a controlled manner and build up a greater degree of confidence.

The one major problem that is not addressed in the current experiment is spatial gradients in the temperature and density (7.39). Indeed in Figs. 7.32 and 7.33, we see some evidence that the laser at long times is heating up a larger volume of under dense plasma and thus biases the electron density. To test that this is the case one needs to design experiments which have both temporal and spatial resolution. This can be achieved by either using "spot" spectroscopy in conjunction with a streak camera (Herbst and Grun, 1981), (7.39) or by using a framing camera (which

could also employ spots). The main point is that now with a proven ability to do time dependent electron density measurements the next advances are possible.

7.6.6 Conclusion

In conclusion the results of recent experiments on laser irradiated disk using an x-ray crystal streak camera show that time dependent electron densities can be determined. The method of using Stark width measurements on suitable target materials has some major advantages.

- Numerous elements can be used to give results on n_e over a wide range of n_e and T_e .
- The results are independent of ionization balance models unlike line ratios.
- The widths are a sensitive function of n_e . Thus one can estimate errors at less than 20%.

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7.7 Studying Electron Distributions Using the Time-Resolved Free-Bound Spectra From Coronal Plasmas

7.7.1 Introduction

Laser light absorption (7.47) (7.48) (7.49) by inverse bremsstrahlung remains an important process by which energy can be efficiently delivered to inertial confinement fusion targets. Light absorption by other physical processes such as resonance absorption (7.50) (7.51) or ion-acoustic turbulence (7.49) (7.52) (7.53) are also important but can lead to production of high energy electrons which preheat and therefore degrade the performance of an ablative implosion. The trend in present day fusion targets is toward higher Z plasmas produced by low laser intensities, short laser wavelengths and long laser pulses. These conditions are of course optimum for absorption by inverse bremsstrahlung. However for

materials with $Z \gg 1$ and at even moderate laser intensities, absorption by inverse bremsstrahlung is expected to be reduced due to non-linear effects. In fact, Langdon (7.47) has calculated that when $Z(v_o/v_{te})^2 > 1$, inverse bremsstrahlung results in a non-Maxwellian velocity distribution weighted toward higher electron energies than are optimum for the absorption process. In this work we investigate the electron distribution by measuring the energy spectrum of aluminum K-shell recombination radiation. We demonstrate that the shape of the recombination spectra can be used to determine the electron distribution function. To optimize absorption by inverse bremsstrahlung, the experiments are performed with green light ($\lambda_L = 0.532 \mu\text{m}$) and with fairly long pulse lengths ($\tau_L \sim 900$ psec). The $Z(v_o/v_{te})^2 \geq 1$ criterion is predicted to be satisfied by performing the experiments with aluminum ($Z = 13$) and $I_L \sim 10^{16} \text{W/cm}^2$.

7.7.2 Definition of Experiment

Laser light absorption by inverse bremsstrahlung is due to classical electron-ion collisions. Langdon (7.47) demonstrated that the time evolution of an initially monoenergetic electron distribution, v_o , under the perturbation of electron-ion collisions -- assuming $Z \gg 1$ and neglecting e^-e^- collisions -- progresses towards a flat (or constant) distribution for $v_{te} < v_o$. Here, v_{te} is the electron thermal velocity $\frac{KT}{m}$ and v_o is the electron "quiver" velocity -- corresponding to light absorption at an electron ion collision rate equal to the laser frequency. The phenomenon occurs if the electrons gain energy from the laser faster than they can thermalize (through e^-e^- collisions) or if

$$Z \frac{e E_L^2}{m \omega_L v_{te}} > 1 \quad (1)$$

where E_L is the laser energy and ω_L is the laser frequency.

In Figure 7.34 we show Langdon's results for the electron velocity distribution function, $f(v)$, of a Maxwellian and a strong non-Maxwellian

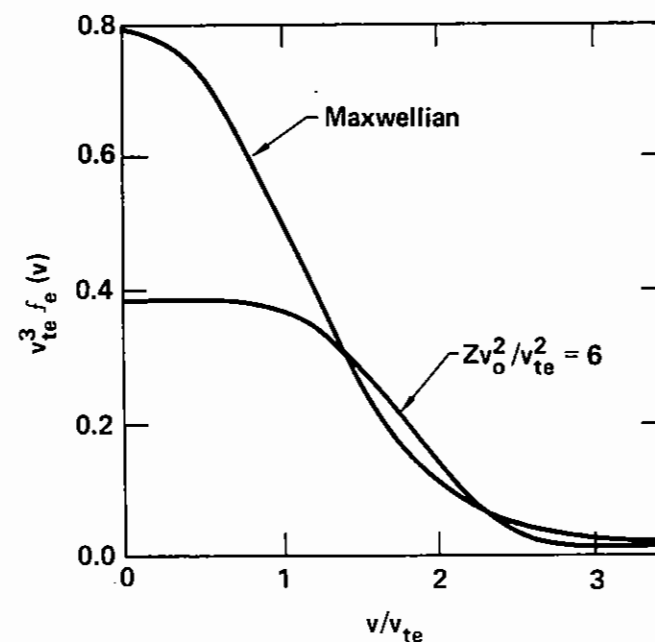


Fig. 7.34 Non-Maxwellian electron distribution function produced by inverse bremsstrahlung absorption when $Z(v_o/v_{te})^2 > 1$. The two curves show distributions when $Z(v_o/v_{te})^2 \leq 1$ (Maxwellian) and $Z(v_o/v_{te})^2 = 6$.

$(Z(v_o/v_t)^2 = 6)$. For the measurements discussed here, $Z(v_o/v_t)^2 \sim 7$ since $(v_o/v_t)^2 = 4 \times 10^{-16} I \lambda^2 / T_e$, where I (W/cm²), λ (microns), T_e (keV) assuming $Z = 13$ (fully ionized), $T_e \sim 1$ keV, $I \sim 0.5 - 1.0 \times 10^{16}$ W/cm² and $\lambda_t = 0.532$ μ m.

The derivation of the emission rate for the free-bound continuum of a hydrogenic ion is obtained by noting that photoionization and radiative recombination are inverse processes, and are in detailed balance. Since we can derive the free-bound emission rate per ion and electron having velocity v , then the non Maxwellian (non-thermal) emission can be determined by using the appropriate non-Maxwellian distribution function.

First the absorption coefficient for frequency ν from hydrogenic ion in state n' .

$$\alpha_\nu = \frac{64 e^{10} \pi m^2}{3 \sqrt{3} h e} \frac{1}{(n')^5 \nu} \quad (2)$$

Defining $F(\nu)$ as the probability of there being an electron with velocity

v then we know that

$$n_I n_e(\nu) F(\nu) v \, dv \, dt \quad (3)$$

is the number of "recaptures" and ν is related to velocity by

$$h\nu = E_{\text{IONIZATION}} + 1/2 m v^2, \quad (4)$$

i.e., by energy conservation. Therefore the emission per unit frequency per unit is

$$n_\nu = h\nu n_I n_e(\nu) F(\nu) v \, dv = \left(\frac{h^2 \nu}{m}\right) n_I n_e(\nu) F(\nu) \, dv \quad (5)$$

To relate $F(\nu)$ to the absorption cross-section α_ν of Eq.2, we use the continuum analog of the Einstein-Milne relations and find

$$\alpha_\nu = \frac{4\pi c^2 m^2 \nu^2}{h^2 \nu^2} \frac{g_1}{g_0} F(\nu) \quad (6)$$

here we used g_1, g_0 as the degeneracy of higher and lower ionization stage. Substitution leads to

$$n_\nu = Z^4 n_I n_e(\nu) \frac{e^{10} 16 \pi^3}{3 \sqrt{3} m^2 c^3 h^2} \frac{g_0}{g_1} \frac{1}{\nu^2} \quad (7)$$

we thus have for the recombination continuum

$$n_\nu = \frac{n_e(\nu)}{\nu^2} \cdot A \quad (8)$$

where A contains the constants that do not depend on the velocity distribution. The emission can now be characterized by specifying the velocity distribution;

$$n_e(\nu) = n_0 f(\nu) \quad (9)$$

for two cases:

a) Maxwell Boltzmann (M.B.),

$$\frac{f(\nu)}{\nu^2} = 4\pi \left(\frac{m}{2\pi k T}\right)^{3/2} \exp \left\{ -\left(\frac{m\nu^2}{2kT}\right) \right\} \quad (10a)$$

b) a flat topped (F.T.) distribution which represents the asymptotic self-similar solution relevant to the Langdon effect,

$$\left(\frac{1}{\nu_0^3}\right) \exp \left\{ -\left(\frac{\nu}{\nu_0}\right)^5 \right\}$$

This distribution can be approximately represented as:

$$\begin{aligned} \frac{f(\nu)}{\nu^2} &= \frac{3}{\nu_0^3} \quad \nu < \nu_0 \\ &= 0 \quad \nu > \nu_0 \end{aligned} \quad (10b)$$

Note that the distribution is a speed distribution and that v_0 is the velocity at some time when the original distribution is no longer part of the evolution. These two distributions, together with Eq. 8, can be rearranged to show that the logarithm of the emissivity as a function of frequency is given by:

$$\text{M.B.: } \log(n_v) \propto - (h\nu/kT) \quad (11a)$$

$$\begin{aligned} \text{F.T.: } \log(n_v) &\propto \text{constant for } h\nu > E_I + 1/2 m v_0^2 \quad (11b) \\ &\propto 0 \text{ for } h\nu > E_I + \frac{1}{2} m v_0^2 \end{aligned}$$

Thus one would expect a constant slope for the F.T. distribution with a rapid decrease above a certain photon energy.

7.7.3 Experimental Method

A schematic of our experiment is shown in Figure 7.35. In essence, we measured the time evolution of the K-shell free bound spectrum that is emitted from the coronal plasma produced by a laser-irradiated aluminium disk. We utilised the single-beam Vulcan laser at the Rutherford Appleton Laboratory's Central Laser Facility. Laser output parameters were nominally ~ 800 ps pulsewidth and $\lambda_L = 0.53 \mu\text{m}$ wavelength. The laser was focused to a minimum spot size of $\sim 50 \mu\text{m}$ diameter which produced intensities of $I_L \sim 0.5 - 1.0 \times 10^{16} \text{ W/cm}^2$, the variation owing to fluctuations in laser energy output (15-30 joules). The surface of the disk was oriented perpendicular to the direction of the laser beam. A streaked X-ray spectrograph (7.54), which utilized a RAP crystal to diffract photons into an X-ray streak camera, was positioned at 90° relative to the incident laser beam. The geometry of the spectrograph plus the limited width of the streak-camera cathode (2 cm) restricted the energy coverage to the interval 2.2 keV - 3.0 keV. This interval is adjustable in terms of centroid energy but not width. For more details concerning this instrument see reference (7.54). This energy interval was chosen to cover the Al XIII recombination edge and higher energies. The relative efficiency of the streak camera spectrograph was obtained by normalizing to a spectrum obtained simultaneously using a conventional time-integrated spectrograph. This normalization requires the integration

of the time-resolved spectrum over the measured time interval (0 to 2.5 ns).

To ensure that the spectrograph measured only X-rays emitted from the coronal plasma, the viewing region is restricted using a mask located 25 μm or more from the disk surface; as can be seen in Figure 7.35. The mask is made of Ti that is 25 μm thick and therefore opaque to any X-rays of interest. The mask is placed adjacent to the target (at $\sim 50 \mu\text{m}$ from the edge of the disk). Density profiles obtained by interferometry (7.55) for laser produced plasmas generated by similar conditions indicate that with the plasma masked as indicated, we are only looking at the corona.

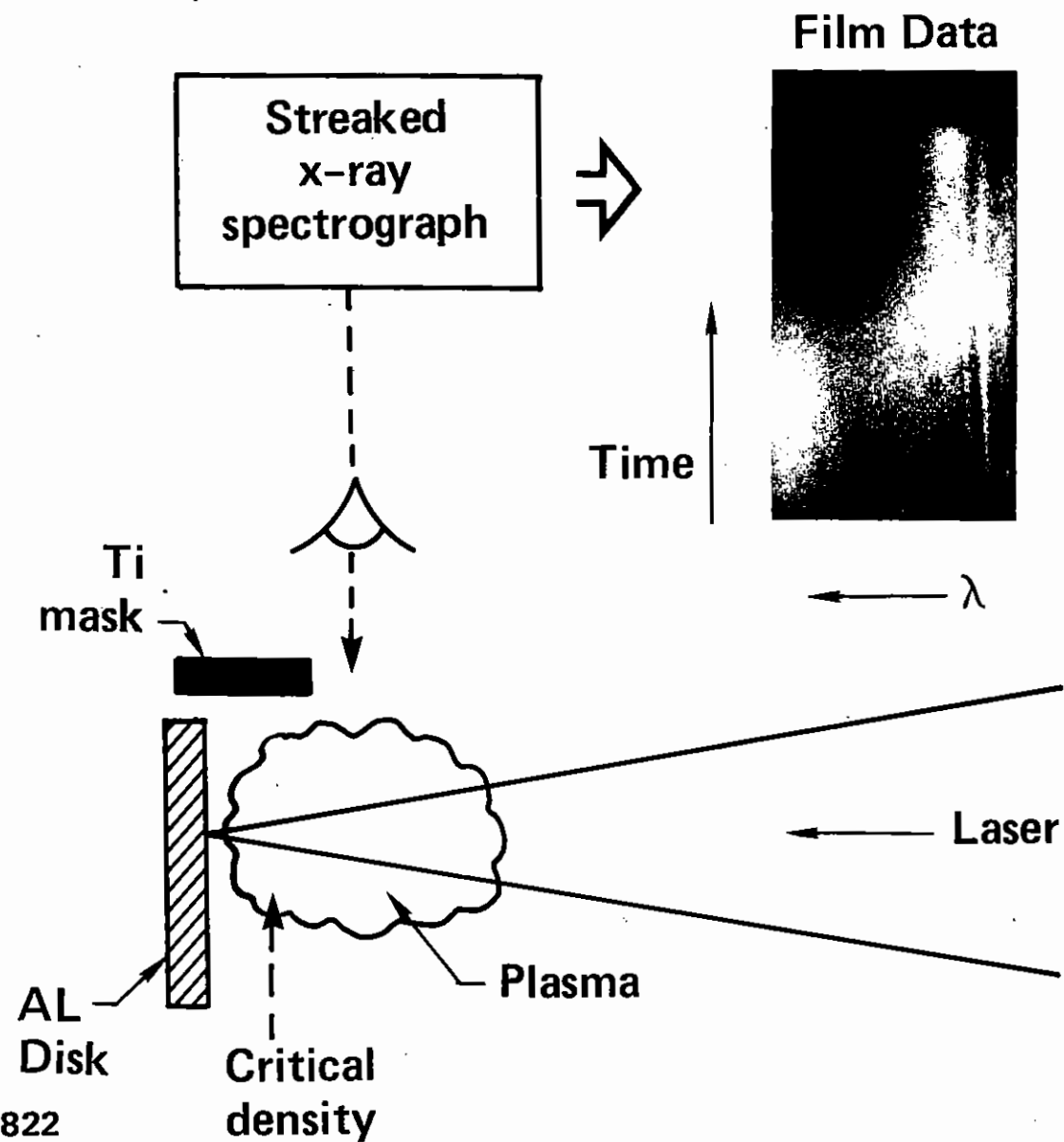
In Figure 7.35 is a picture of the time-resolved spectrum. The signal/noise ratio is sufficient to observe the free-bound emission spectrum for a time interval ~ 2 ns. This interval of emission far exceeds the laser pulselength (FWHM).

Figure 7.36 represents a 2-D image of the film produced by a storage of film density vs position on a VAX 11/780 computer system. In figures 7.37(a)-(d) and 7.38(a)-(c) we show spectra corresponding to different time intervals corresponding to when the laser is on and then off. As marked in Figure 7.37 these spectra are obtained by averaging the energy spectrum for a given time interval.

Figure 7.39 represents a typical spectrum obtained with a time-integrated spectrograph utilized in this same set of measurements. The illustrated spectral regions have similar shape in both time-resolved and time-integrated spectrographs only if the time resolved data is integrated over the entire observation interval. Otherwise the spectra exhibit marked differences for a given time interval. If you wish to determine electron temperature in a laser-produced plasma, this phenomenon should illustrate the futility of using the time-integrated spectrum of freebound emission.

7.7.4 Discussion

The spectra shown in Figure 7.37(a)-(d) represent the emission observed



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Fig. 7.35 Schematic of Experiment.

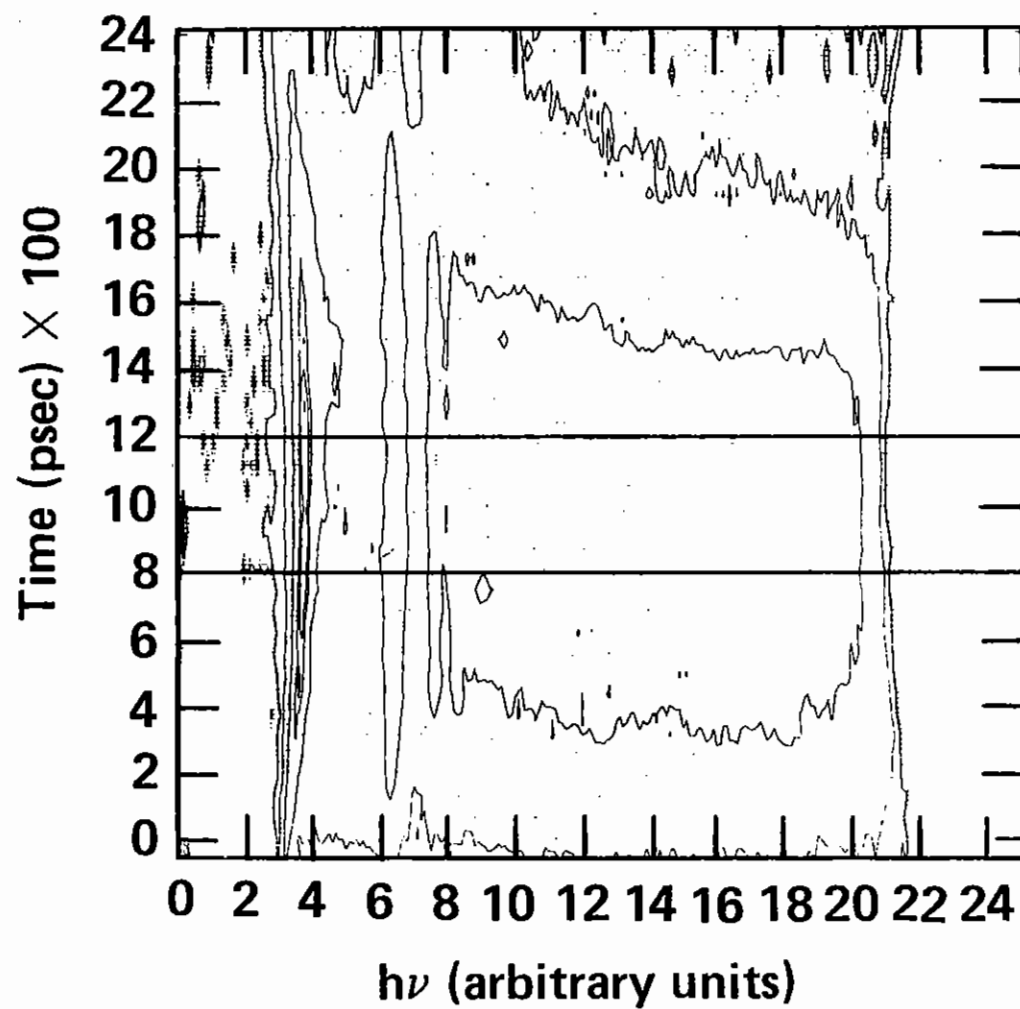


Fig. 7.36 2-D representation of digitized film data. Shaded region illustrates time interval that can be arbitrarily set to take a snapshot of x-ray spectrum.

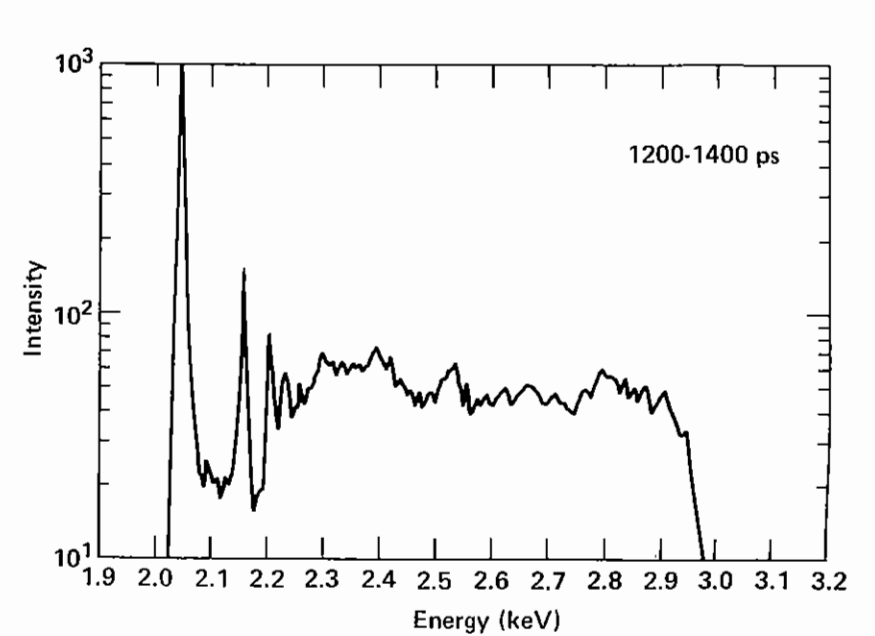
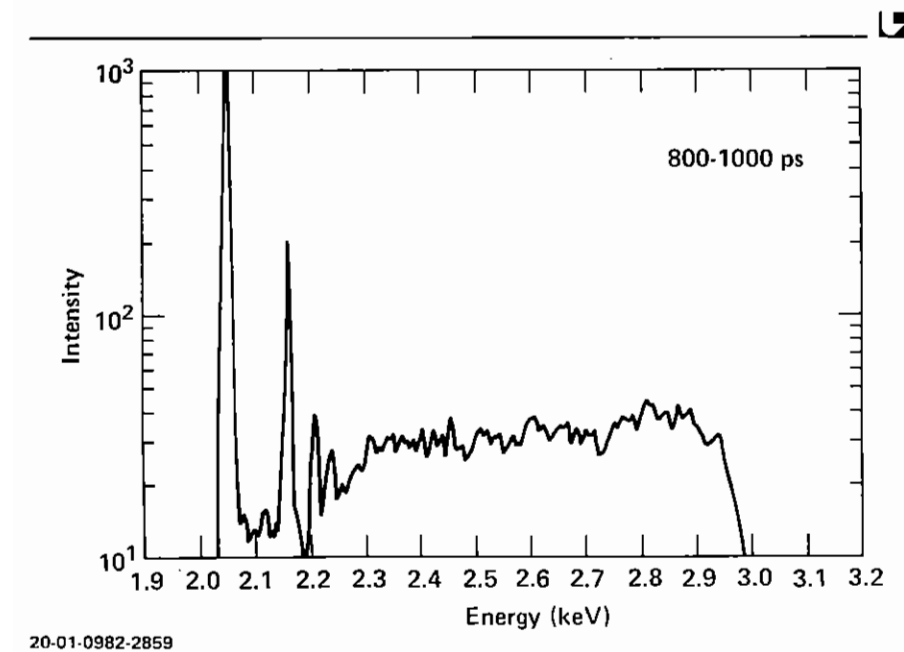
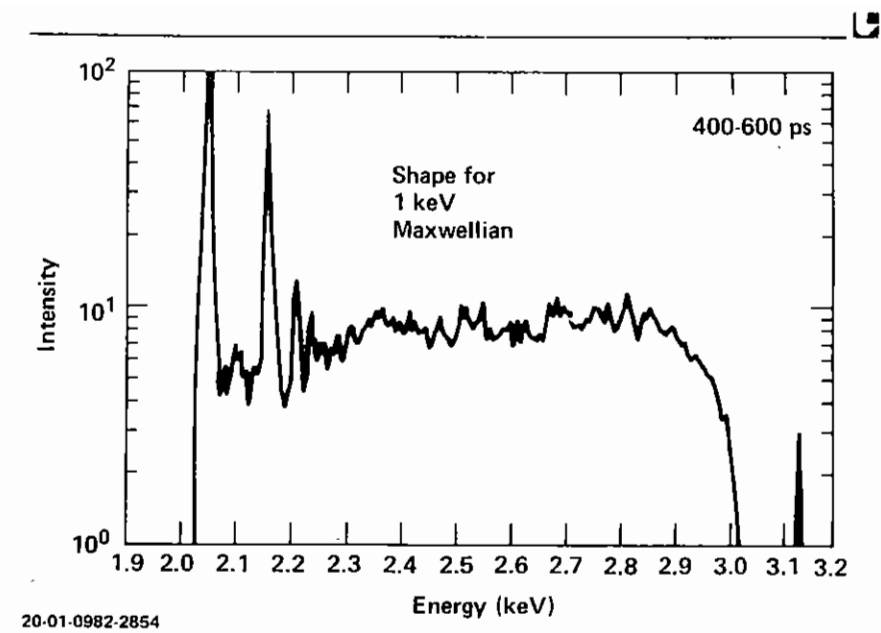
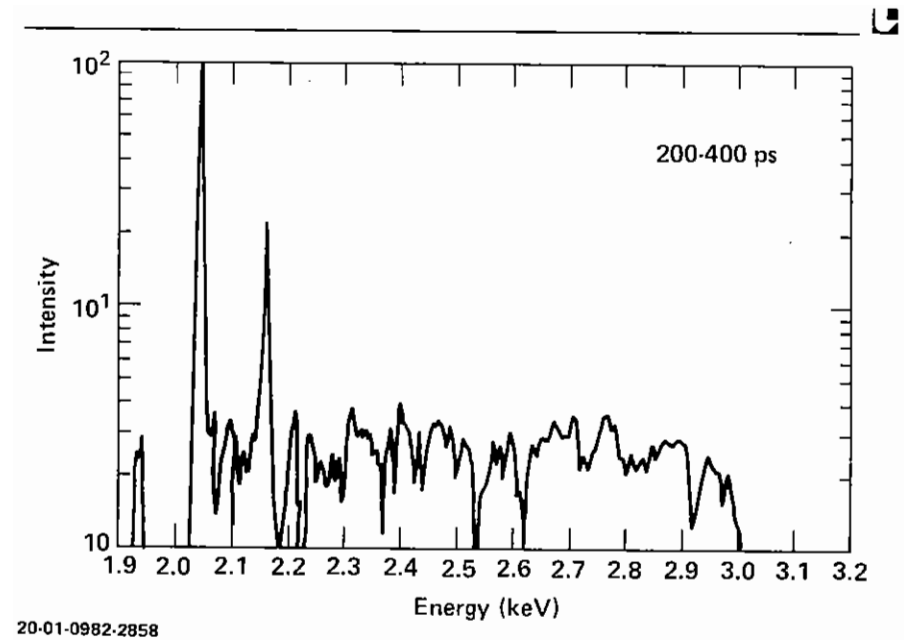


Fig. 7.17 Snapshot of Al K α -ray spectrum obtained over noted time intervals which corresponds when laser is heating plasma.

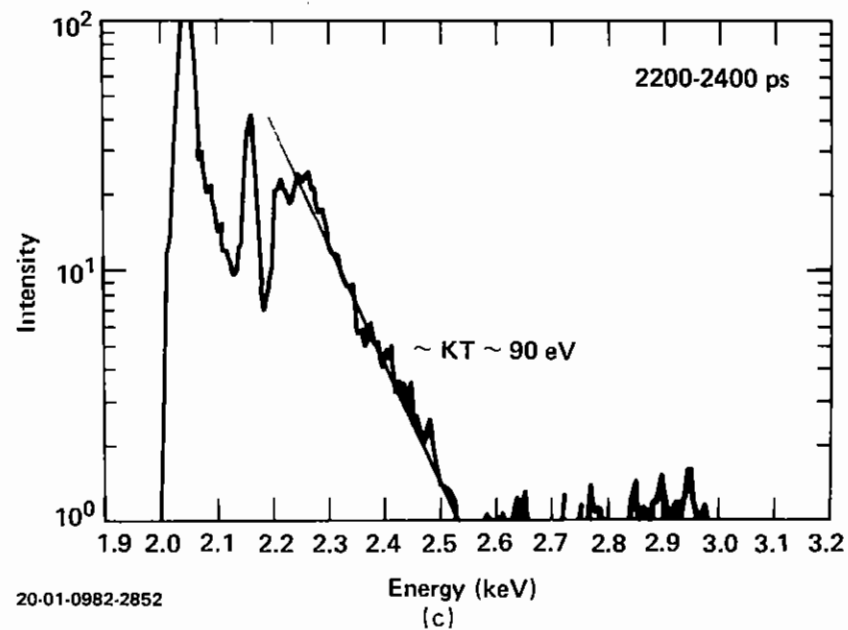
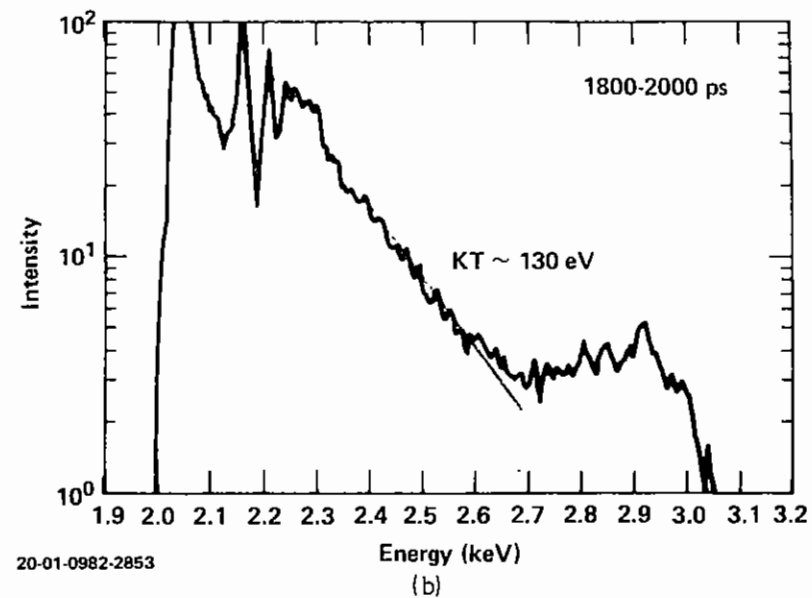
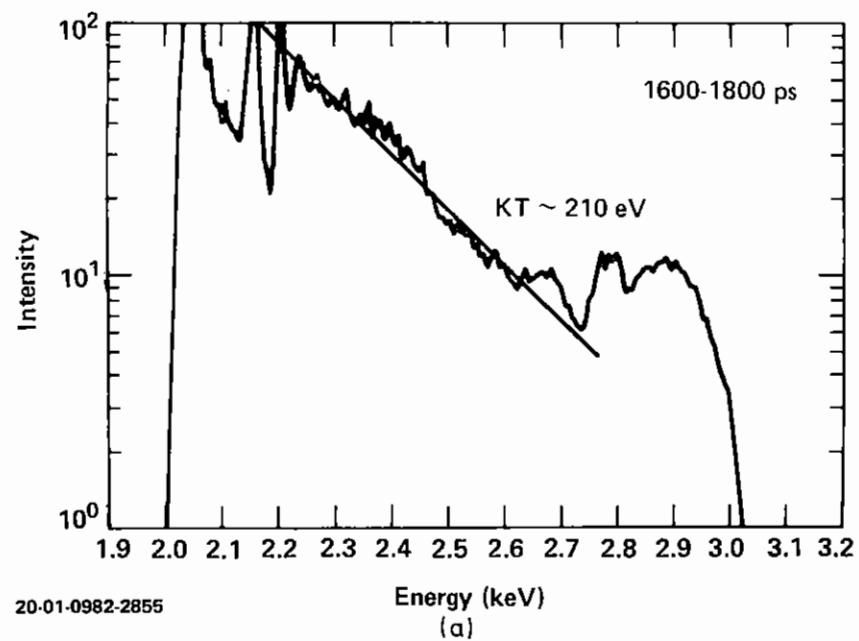


Fig. 7.38 Snapshots of Al K X-ray spectrum obtained over noted time intervals which correspond to after laser heating is shut off.

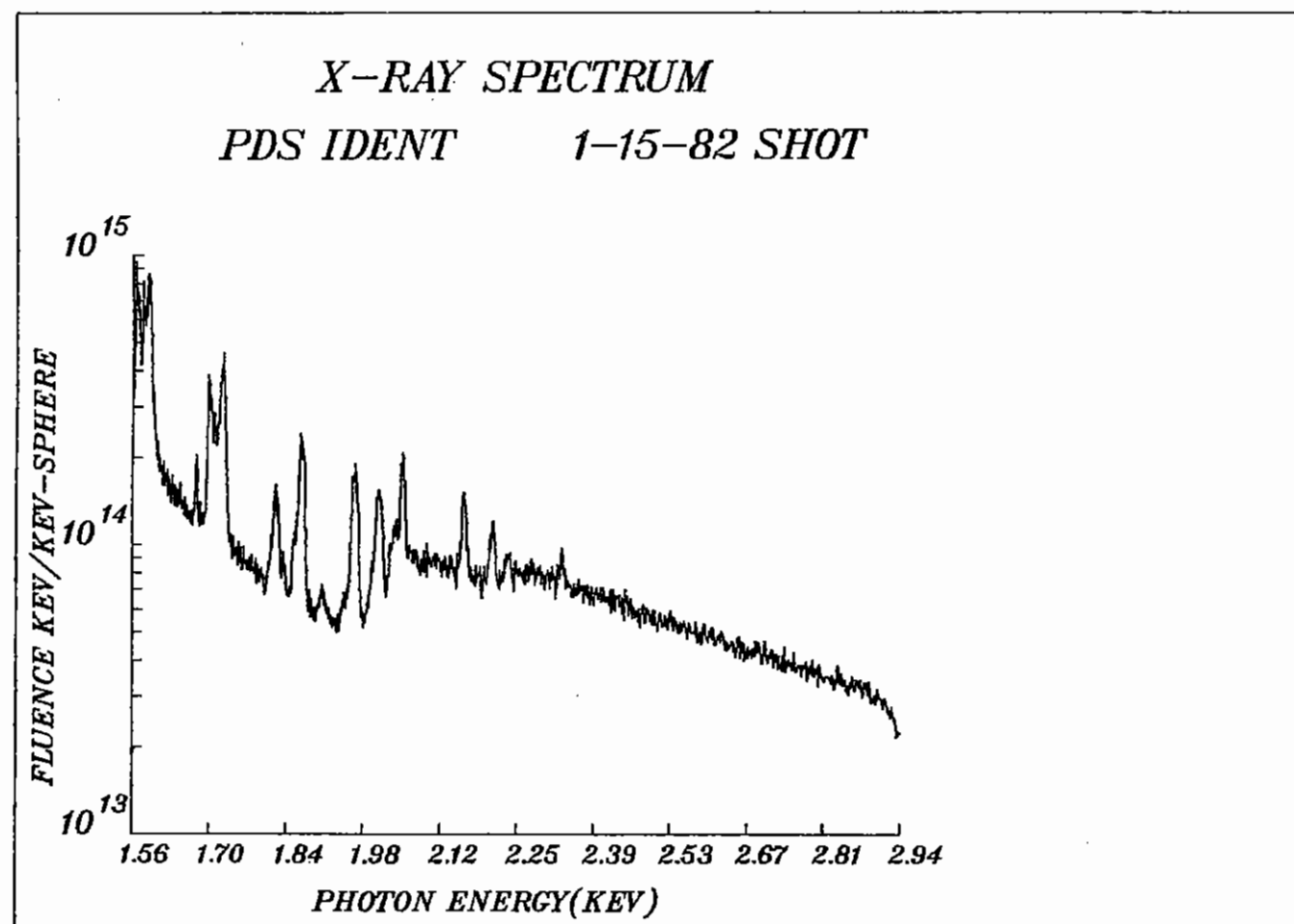


Fig. 7.39 Al K x-ray spectrum obtained with conventional time-integrating spectrograph.

during the time intervals 200-400 ps, 400-600 ps, 800-1000 ps, and 1200-1400 ps, respectively. These times roughly correspond to those during which the laser is heating the plasma (recall the pulselength is ~ 800 ps FWHM). The spectra produced during the first three time intervals are nearly identical except for relative emission levels, the 400-600 ps interval being the most intense. These spectra are characterized by nearly constant intensity emission from approximately 2.2 to 2.9 keV. The rolloff in intensity above 2.9 keV is an effect due to the finite width of the streak camera cathode. On Figure 7.37(b) we have drawn the spectral shape which would be expected from a 1 keV Maxwellian. The fact that the observed spectrum is much flatter than expected for the Maxwellian indicates that, either KT is much higher ($KT > 6$ keV would appear flat over this narrow energy interval) or the spectrum is non-Maxwellian. The spectrum does not drop rapidly for energies larger than the threshold energy. This may be due to an electron distribution $f(v)$ which is neither constant nor Maxwellian. The 1200-1400 ps interval spectrum begins to illustrate some slope towards lower intensity at higher photon energies. This should indicate some cooling of the plasma late in the laser pulse.

Figure 7.38(a)-(c) illustrates a method to obtain the time rate of change (or cooling rate) of the plasma temperature, provided the laser is off. Presuming that late in the lifetime of the plasma we can characterize $f(v)$ as a Maxwellian, then it should be valid to use the slope of the freebound emission to determine electron temperature. Under these circumstances it is easy to show that the plasma cools from $KT \sim 210$ eV in the 1600-1800 ps interval down to $KT \sim 90$ eV in the 2200-2400 ps time segment.

There appears a "hump" in all three spectra represented in Figures 7.38(a)-(c) which is centred at $h\nu \sim 2.9$ keV. This effect is instrumental, although we have not isolated a cause. If it is real, it represents contributions from electrons having some net velocity relative to the plasma as a whole. This phenomenon would produce a positive shift in the spectrum corresponding to the equivalent velocity of the plasma centre of mass relative to the nonisotropic electron source. A net velocity of free electrons relative to the capturing ions of $v_e \sim 1.5 \times 10^9$ cm/sec would be sufficient to shift the recombination

spectrum by ~ 600 eV. As another possible explanation, we may be observing an isotropic but non-thermal distribution which peaks around ~ 600 eV (ie, nearly monoenergetic). Further experiments with a better characterized detector will be needed to study this "hump".

7.7.5 Conclusions

We observe a flat-topped spectrum only when the laser is on which is indicative of either a non-Maxwellian $f(v)$ or a $KT \gg 1$ keV. Spectral coverage over a larger energy range is needed to resolve this ambiguity. This will require re-design of our time resolved instrument to lower its dispersion. As an aside we note that we can easily measure the cooling rate of the plasma after the laser is off by measuring the time rate of change in the free-bound slope.

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7.8 Time resolved X-ray line diagnostics of laser-produced plasmas

7.8.1 Discussion

We have performed a series of measurements to investigate the usefulness of X-ray line intensities as electron density and temperature diagnostics. Line intensities from highly stripped ions have been used successfully in diagnosing high temperature astrophysical and magnetically confined plasmas (7.56)(7.57). Line intensities of He-like and H-like species from laser-produced plasmas have also been used by various researchers to estimate electron temperature and densities in the coronal plasma region (7.58). In general, these measurements have been compromised due to the relatively dense, transient nature of laser-produced plasmas. Large temperature and density gradients also complicate the analysis in determining a unique temperature and density. In our experiment we have begun studying these effects to assess their importance for predicting temperature and densities of underdense plasmas.

A single beam of $0.532 \mu\text{m}$ light is focused using an $f/1$ lens onto planar

targets of various Z materials. The average intensity on target is $\sim 10^{15}$ W/cm². The target normal is tilted 10° with respect to the beam. The targets are discs of CH, SiO₂ and high Z materials which have been doped with low concentrations of elements whose lines are to be used as diagnostics. Typically, the dopant levels are less than ten atomic per cent; although for some high Z targets larger concentrations are needed. In all cases the dopant levels should not change significantly the Z or hydrodynamics of the plasma.

Two crystal spectrographs measure the X-ray spectra. A spectrograph located at 80° to the disc normal measures the time-integrated spectra. A 50 μ m slit placed in front of the spectrograph spatially resolves the plasma along the incident beam direction. An example of the data is shown in Figure 7.40. In general, line emission extends away from the absorption region characterized by high continuum emission into the coronal region where only line emission is observed. An X-ray crystal streak camera located on the opposite side of the target time resolves the line emission with 80 ps resolution (7.54). Figure 7.41 shows an example of the data. For the lower Z targets the X-ray line emission lasts significantly longer than the continuum emission. For the high Z targets the X-ray lines tend to follow the laser pulse.

The data from both instruments are analysed deriving intensities for the various X-ray lines. Examples of the time-integrated data for Si doped in CH and PbO targets are shown in Figures 7.42 and 7.43. Similar spectra were also obtained with the time-resolved spectrograph.

The model to calculate the line ratios used in the study is a straight forward steady state level population model. The inputs require the temperature, density, and all the appropriate rate constants of interest. In the present case we include all of the ground states up to the Be-like stage and then have excited states as well as ground state for Li-, He-, and H-like ions. In the case of the Li-like and He-like we also include the autoionizing states which form satellites on the resonance transition of the next ion stage.

The rates of primary interest; ie the ionisation rates, the electron collision rates in the He-like ground to n equals 2 complex, and the collision rates between the autoionising state are found in the following references:

Ionisation Rates	Lotz Sampson Sampson & Golden	Z Physics <u>280</u> , 466 (1969) (eg) J Physics B <u>13</u> , 2645 (1980)
Collision Rates in He-like	Vinogradov et al	Sov J Quant El <u>5</u> , 630 (1975)
Collisions between autoionising Li-like like and He-like	Blaha, Jacobs Goett et al	Phys Rev A <u>21</u> , 525 (1980) Atom Data & Nuc Data, <u>25</u> , 185 (1980)

With these rates and standard calculations of the remaining rates, which are less crucial, one can produce the ratios given. The one outstanding problem is to resolve whether the plasma does reach the steady state and for this we ran the problem time dependent using a Lagrangian fluid element as a "test particle" and found that the ions easily come to a steady state when reasonable plasma expansion was used.

Examples of the various ratios are shown in Figure 7.44-7.47. These have been chosen since they are sensitive primarily to a single parameter, either N_e or T_e , and not both simultaneously. Experimental values of Si line ratios from doped CH and PbO are indicated. The results are summarised in Table 7.08. The two targets produce similar results indicating that the plasma parameters are similar in the line emitting region.

The different line ratios do not give a unique plasma temperature and density. The most striking disagreements are with line ratios from different charge states. The He-like resonance to Li-like line indicates a lower T_e than the Lyman to He-like resonance ratio. This may be due to

FIG. 7.40

**EXAMPLE OF A SPACIALLY RESOLVED SPECTRUM
FROM AN Al-DOPED GLASS TARGET**

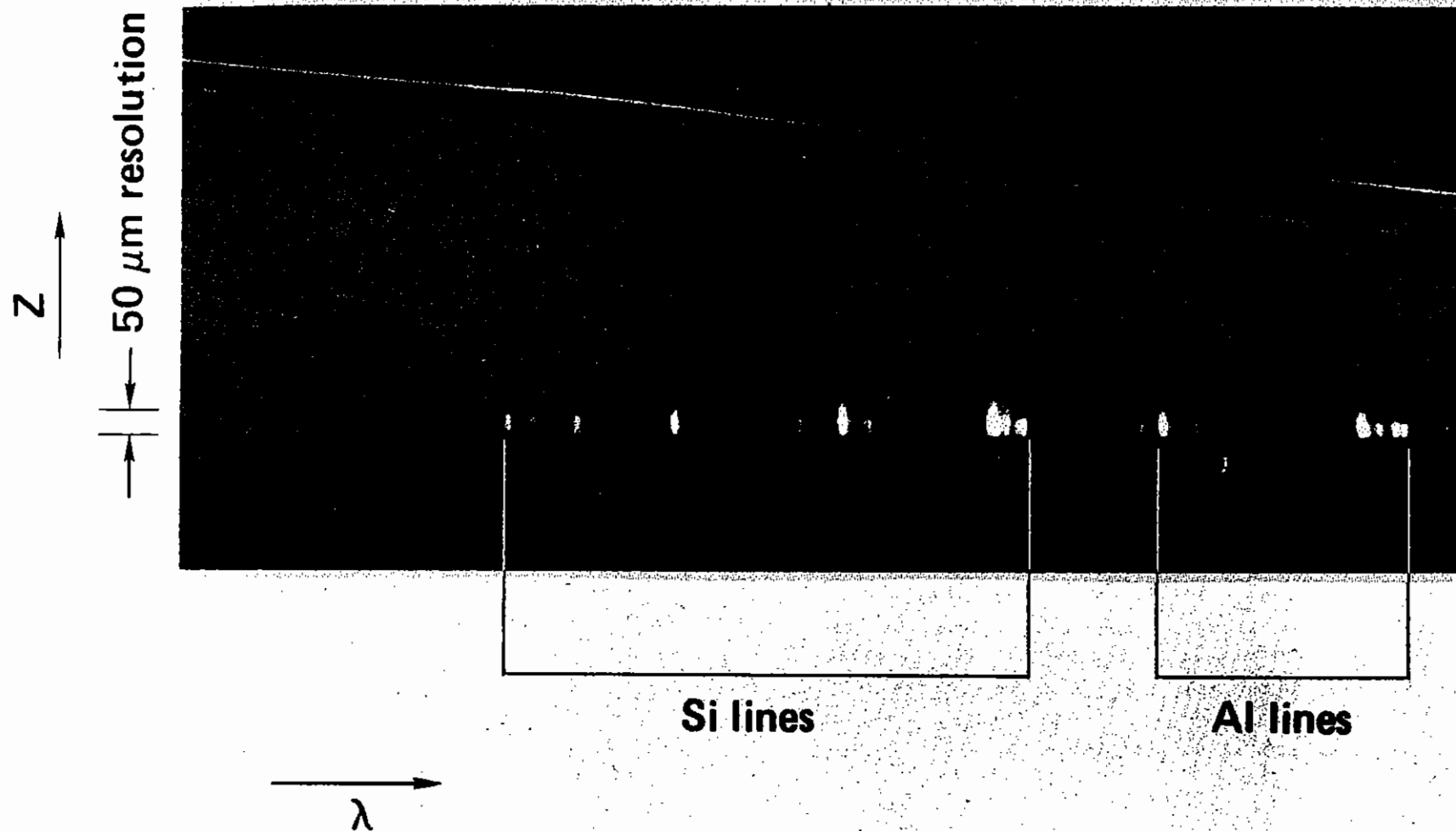
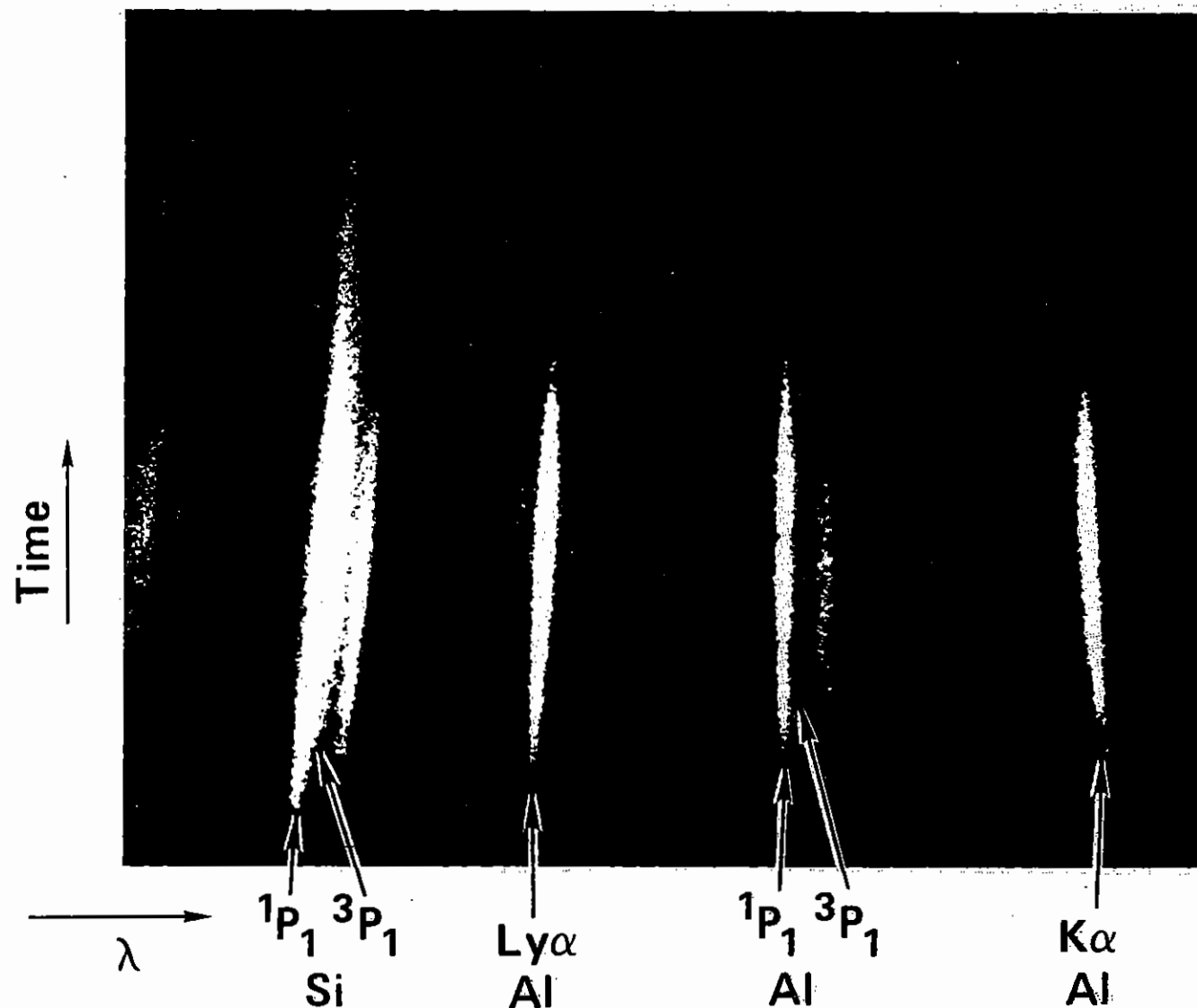


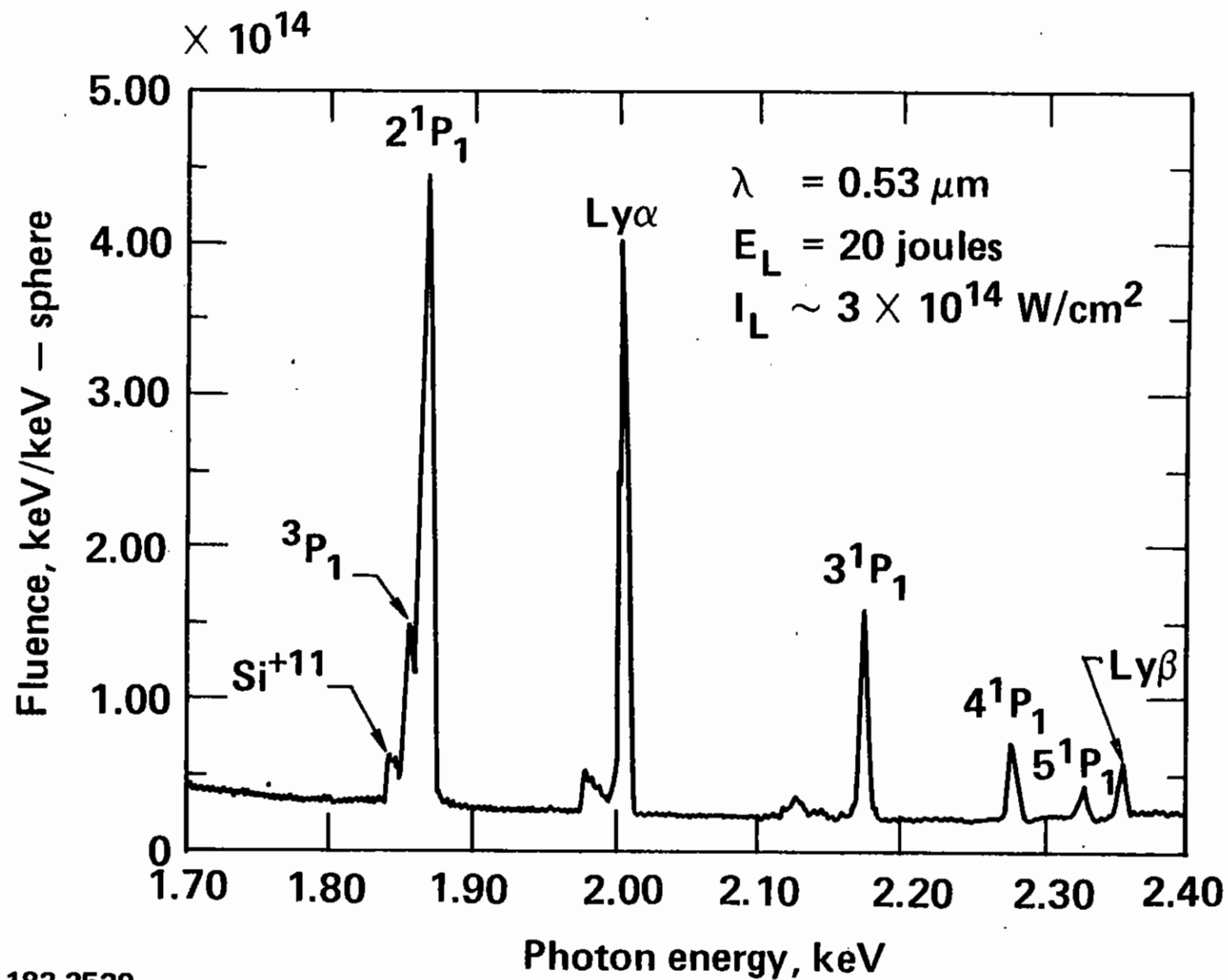
FIG. 7.41
EXAMPLE OF A TIME-RESOLVED SPECTRUM
FROM AN AI DOPED GLASS TARGET



40-90-0482-1657

Shot 1/7/82 #27
10/82

FIG. 7.42.

X-RAY SPECTRUM FROM Si SEEDS IN A CH DISK TARGET

40-01-1182-3539

FIG. 7.43.

X-RAY SPECTRUM FROM Si IN A PbO DISK TARGET

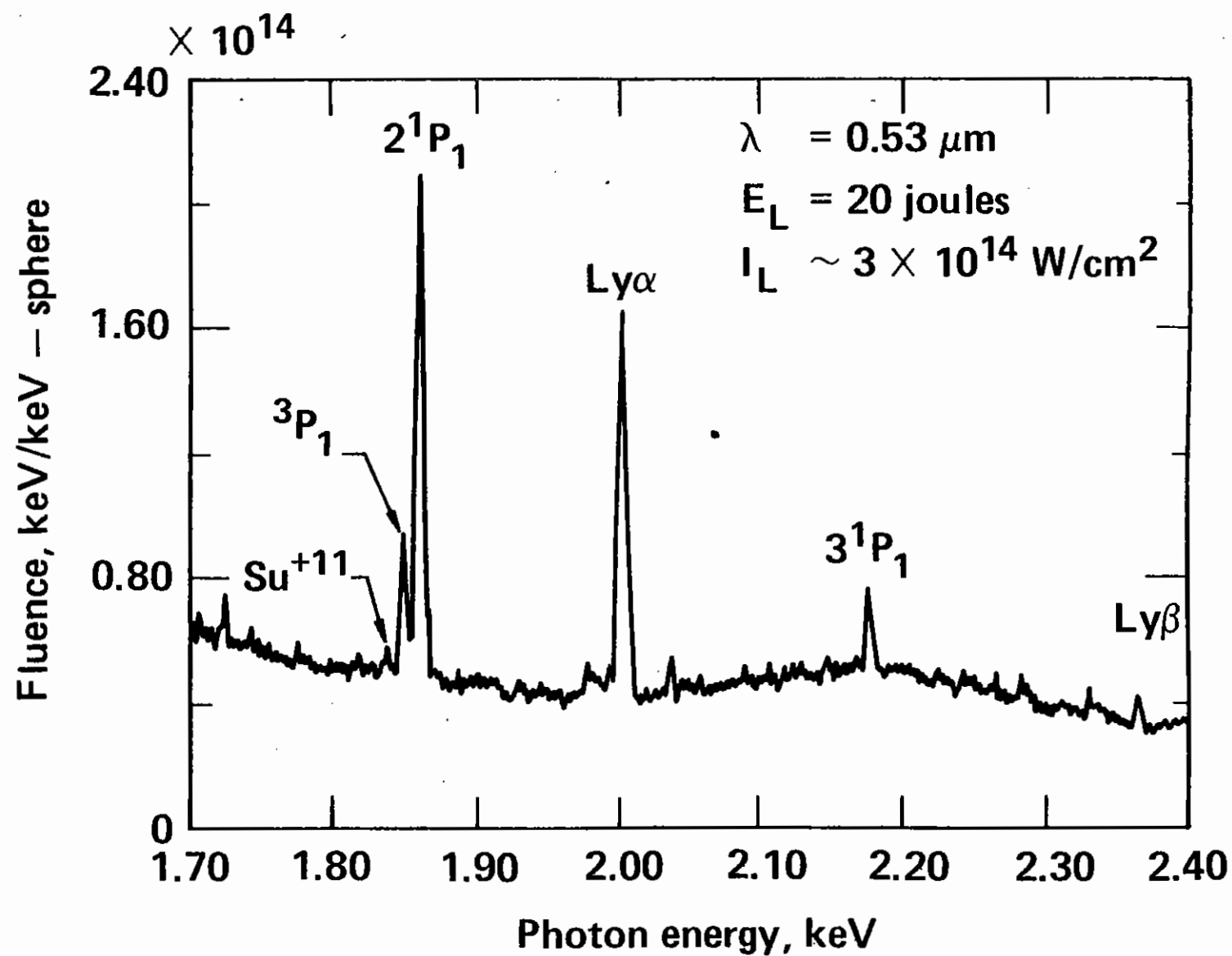


TABLE 7.08

Summary of Plasma Parameters from Si Line Intensities

	$N_e (\text{cm}^{-3})$	
	<u>CH</u>	<u>PbO</u>
$\frac{2^1P_1}{2^3P_1}$	1×10^{20}	1.1×10^{20}
$\frac{3^1P_1}{2^3P_1}$	1×10^{20}	1×10^{20}
$\frac{2^3P_1 + \text{Si}^{+11} (\text{MN+ST})}{\text{Si}^{+11} (\text{ABCD+JKL+QR})}$	$0.3 - 1.1 \times 10^{21}$	$2 - 8 \times 10^{20}$
	<u>Te (keV)</u>	
	<u>CH</u>	<u>PbO</u>
$\frac{\text{Ly}\alpha}{2^1P_1}$	0.83	0.83
$\frac{\text{Ly}\beta}{3^1P_1}$	0.85	1.05
$\frac{2^1P_1}{\text{Si}^{+11} (\text{ABCD+JKL+QR})}$	0.40	0.53

FIG. 7 44

RATIO OF INTERCOMBINATION TO LI-LIKE SATELLITES
IN Si CAN BE USED AS A DENSITY DIAGNOSTIC

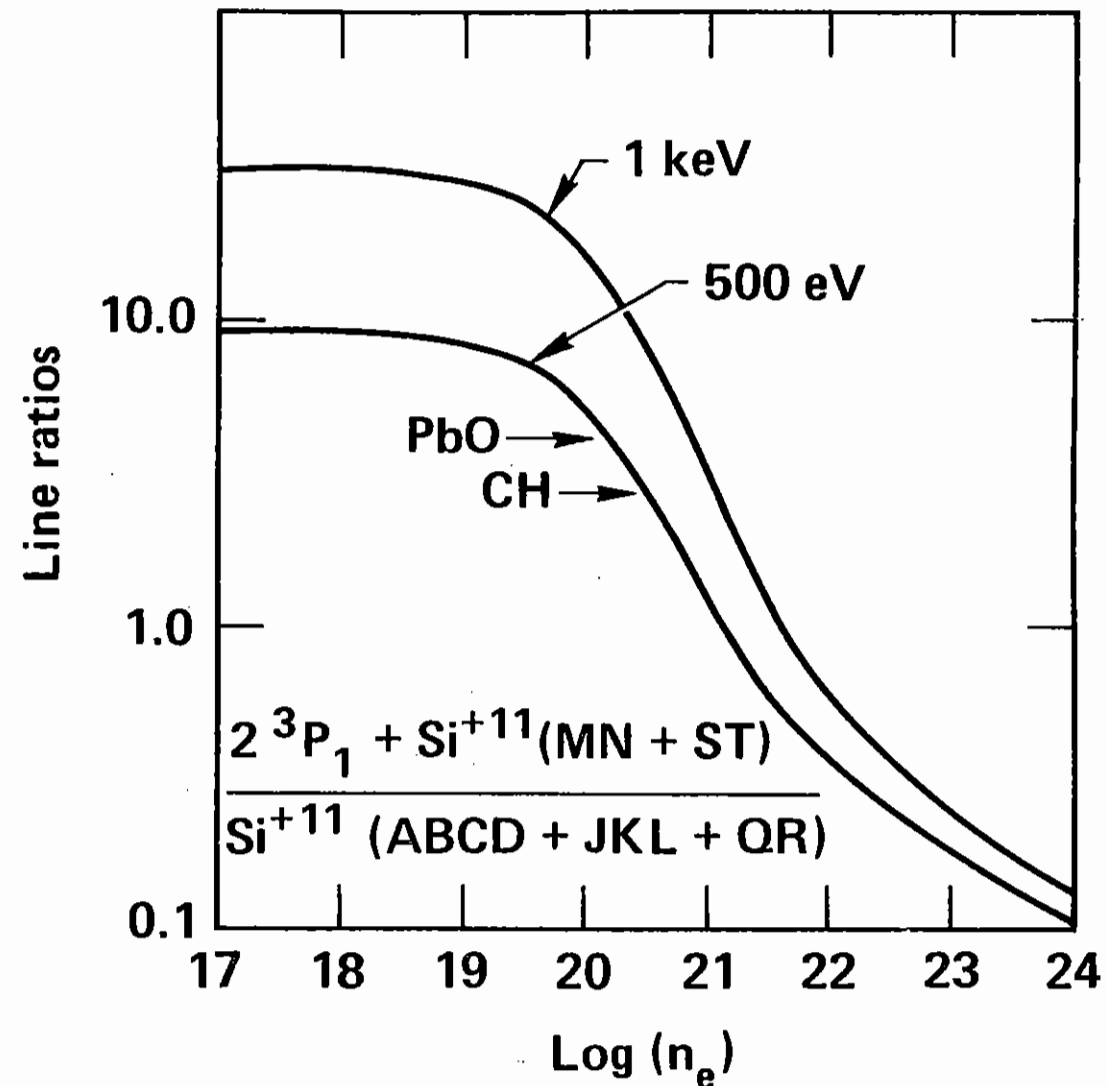


FIG. 7.45.

RATIO OF He-LIKE RESONANCE TO INTERCOMBINATION LINES IN Si CAN BE USED AS DENSITY DIAGNOSTICS

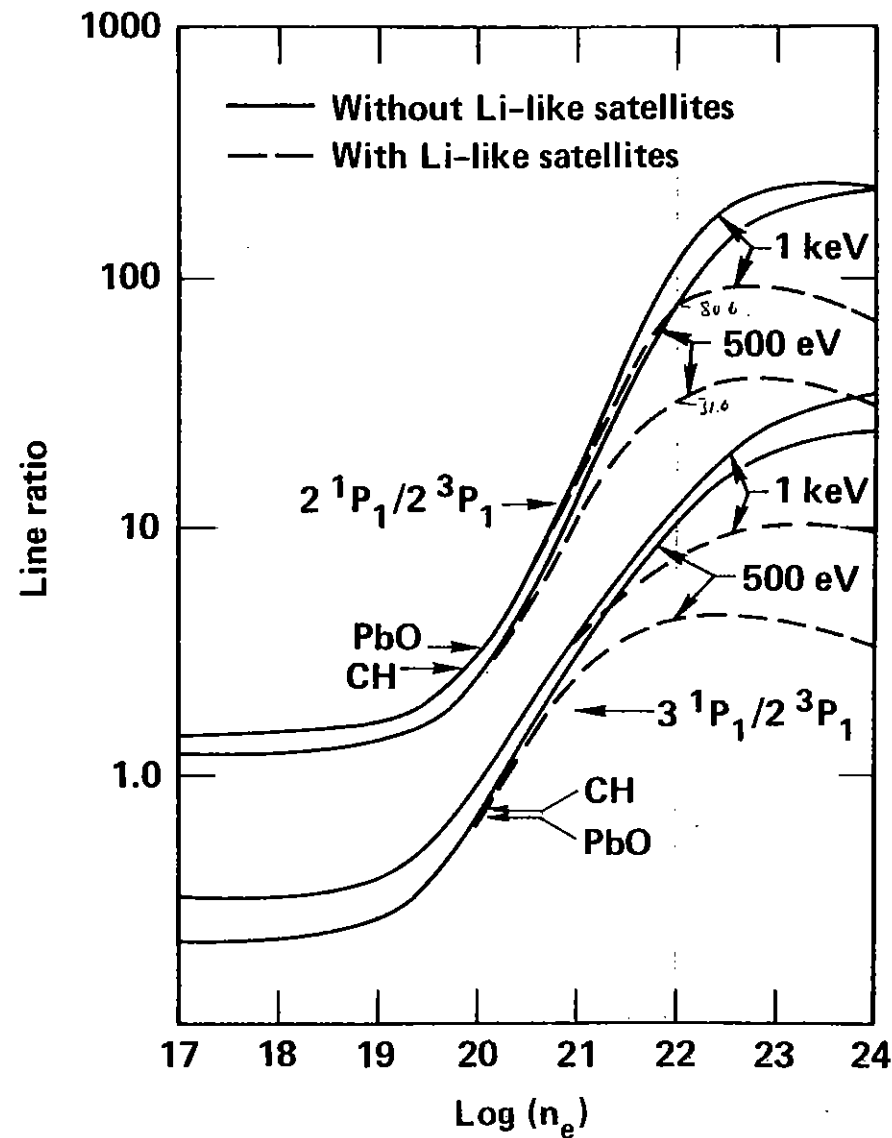


FIG.7.46.

RATIOS OF LYMAN TO HE-LIKE LINES FROM Si CAN BE USED TO ESTIMATE THE LASER PLASMA CORONAL TEMPERATURES

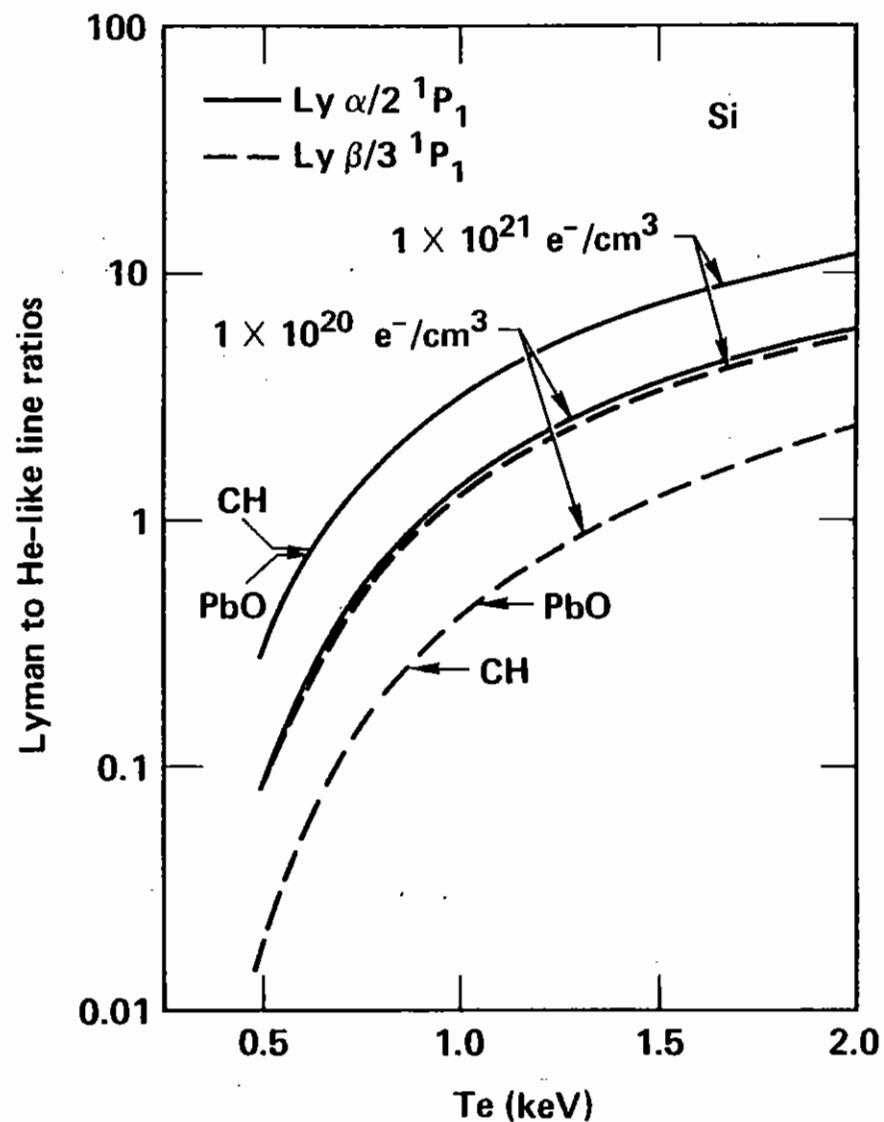
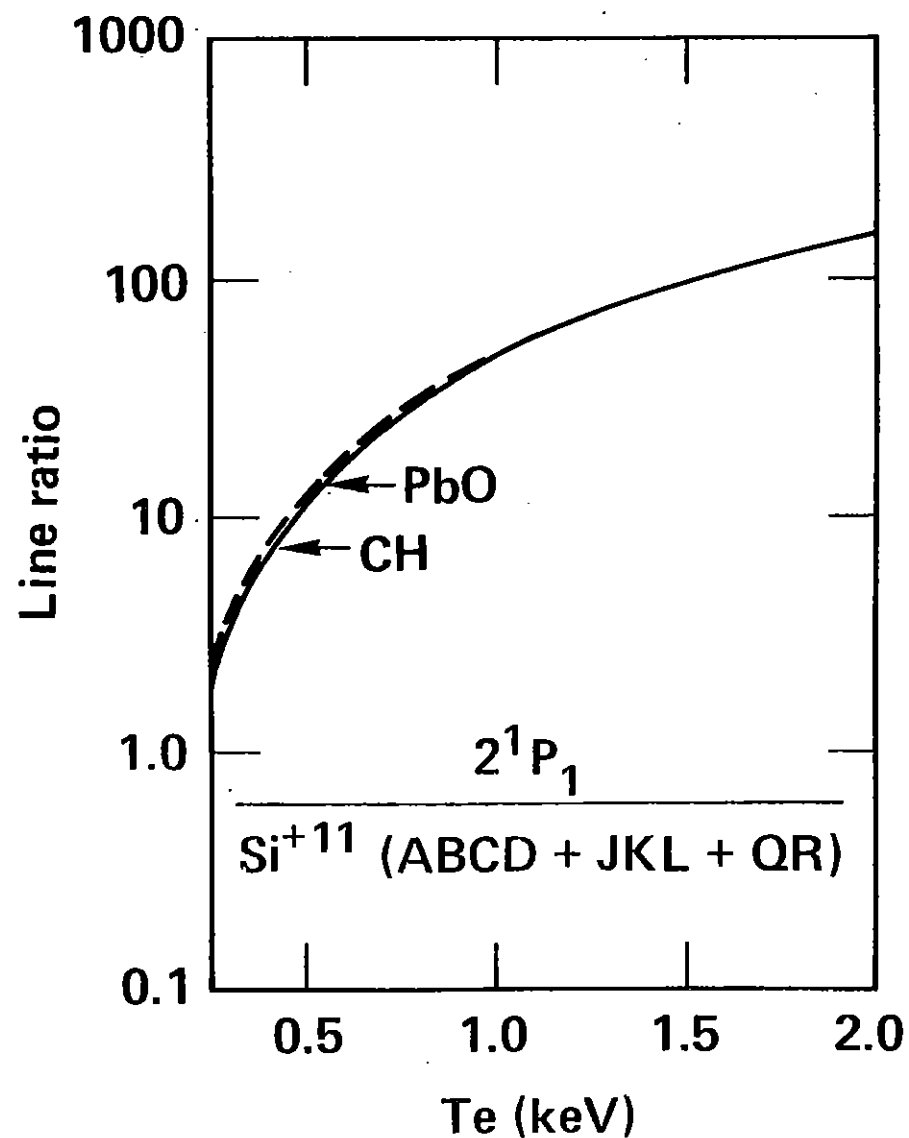


FIG. 7.47

RATIO OF He-LIKE RESONANCE TO Li-LIKE SATELLITES IN Si
CAN BE USED AS A CORONAL TEMPERATURE DIAGNOSTIC



the spatial gradient along the laser beam. Even though the spectrum is spatially-resolved it is averaged over about 50 μm so that large gradients may exist. We plan to use hydrodynamic code predictions coupled with the line ratio code to investigate this effect. We are also investigating ionisation equilibrium and their effects on the line ratios.

An example of the time-resolved data is shown in Figure 7.48. The line ratio of the Si Ly α and He-like 2s¹P₁ line is plotted as a function of time. The time history of the lines are very similar for the two cases.

Finally, various dopants have been seeded into CH targets and their X-ray spectra have been measured. The ¹P₁/³P₁ line ratios are plotted in the Figure 7.49 comparing with predicted line ratios from Boiko et al (7.59). For the seed targets the line ratios indicate emission is from near critical density implying the seeds are sampling the same region of the target.

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7.9 EXAFS Spectroscopy

7.9.1 Introduction

In the EXAFS technique, the fine structure which is present on the high energy side of absorption edges can be used to yield information on the spatial structure, bond lengths, co-ordination number, and chemical identity of surrounding atoms. Unlike X-ray diffraction, which also yields structural information, EXAFS does not rely on long range order in the material. It is also element specific, as a spectrum can be obtained which relates to a specific atom in a structure, and thus is a valuable tool for examining multicomponent systems.

The technique involves measuring the X-ray absorption coefficient as a function of energy above the K, or L, absorption edge of an element. The observed modulation, which is due to constructive or destructive interference between the outgoing part of the photo-electron wave function

with that back-scattered from nearby atoms, depends on the wave number of the ejected photo-electron

$$k = \frac{2\pi}{\lambda} = \frac{1}{\hbar} [2m(E - E_{\text{edge}})]^{\frac{1}{2}}$$

where E is the energy in electron volts and E_{edge} is the energy at the absorption edge of the material.

The 'wiggle' structure which is contained in 50 - 1000 eV above the edge may be described by

$$\mu(E) = \mu_0(E) [1 + \chi(E)]$$

where $\mu(E)$ is the measured absorption coefficient, $\mu_0(E)$ is the absorption coefficient for a single atom in vacuum, and $\chi(E)$ is the fractional modulation of the absorption coefficient.

The accepted formula for EXAFS (7.60) is written as

$$\chi(k) = \frac{m}{4\pi\hbar^2k} \sum_j \frac{N_j}{R_j^2} t_j(k) \exp\left[\frac{-2R_j}{\lambda}\right] \sin[2kR_j + \phi_j(k)] \exp(-2k^2\sigma_j^2)$$

in which N_j is the number of atoms in the j^{th} co-ordination shell, R_j is the average radius of that shell, $t_j(k)$ is the magnitude of the backscattering amplitude for scattering from atoms in the j^{th} sphere, λ is the mean free path for inelastic scattering, $\phi_j(k)$ is the phase shift and σ_j is the root-mean-square deviation of the interatomic distance about R_j , which may contain both static and temperature effects.

To extract information concerning the positions of neighbouring atoms, a Fourier transform of $\chi(k)$ over a range from k_{min} to k_{max}

$$\phi_n(r) = \left(\frac{1}{2\pi}\right)^{\frac{1}{2}} \int_{k_{\text{min}}}^{k_{\text{max}}} k^n \chi(k) \exp(2ikr) dk$$

FIG.7.48.

Si LINE RATIOS CAN BE USED TO DERIVE A TIME RESOLVED CORONAL TEMPERATURE

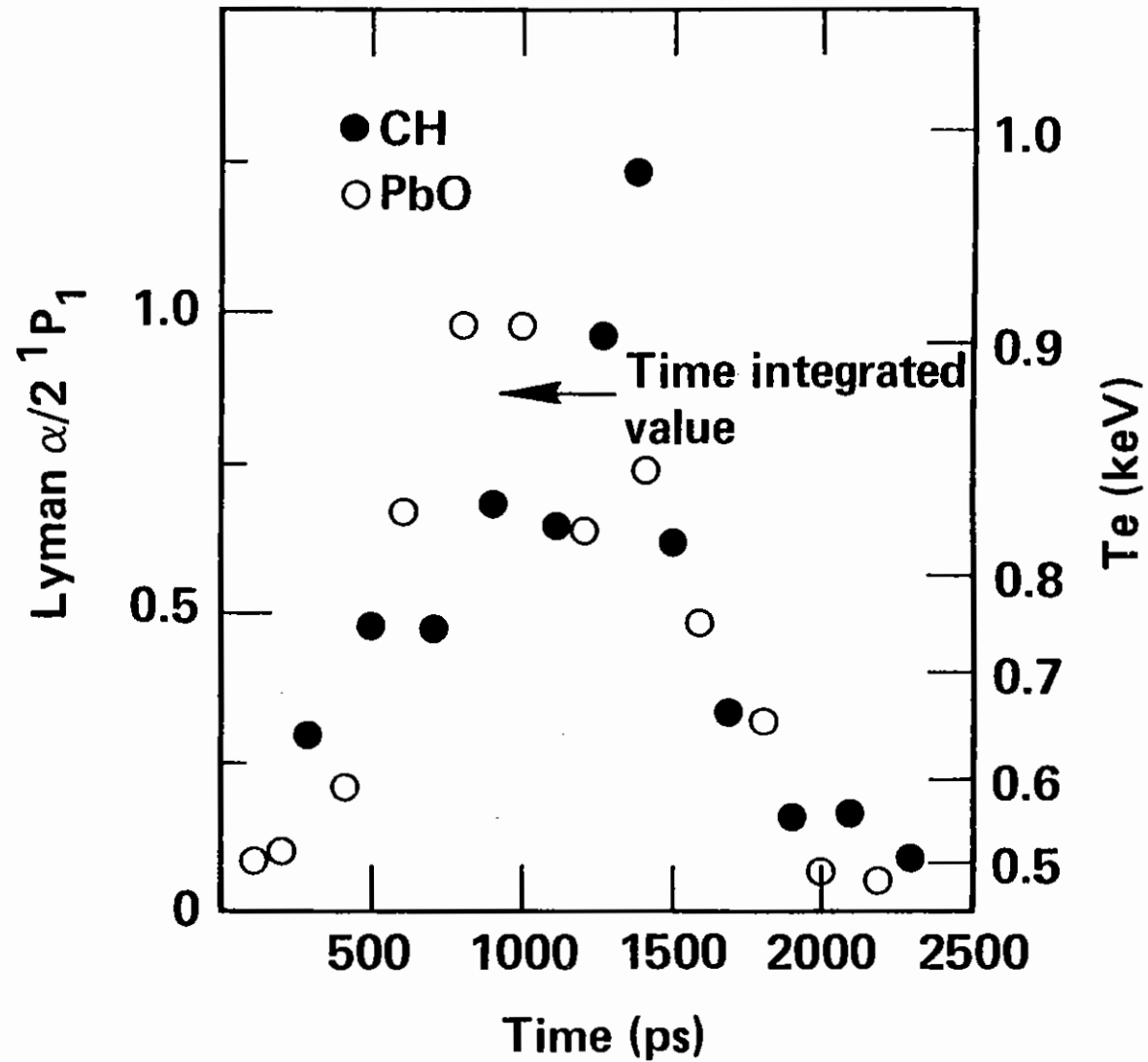
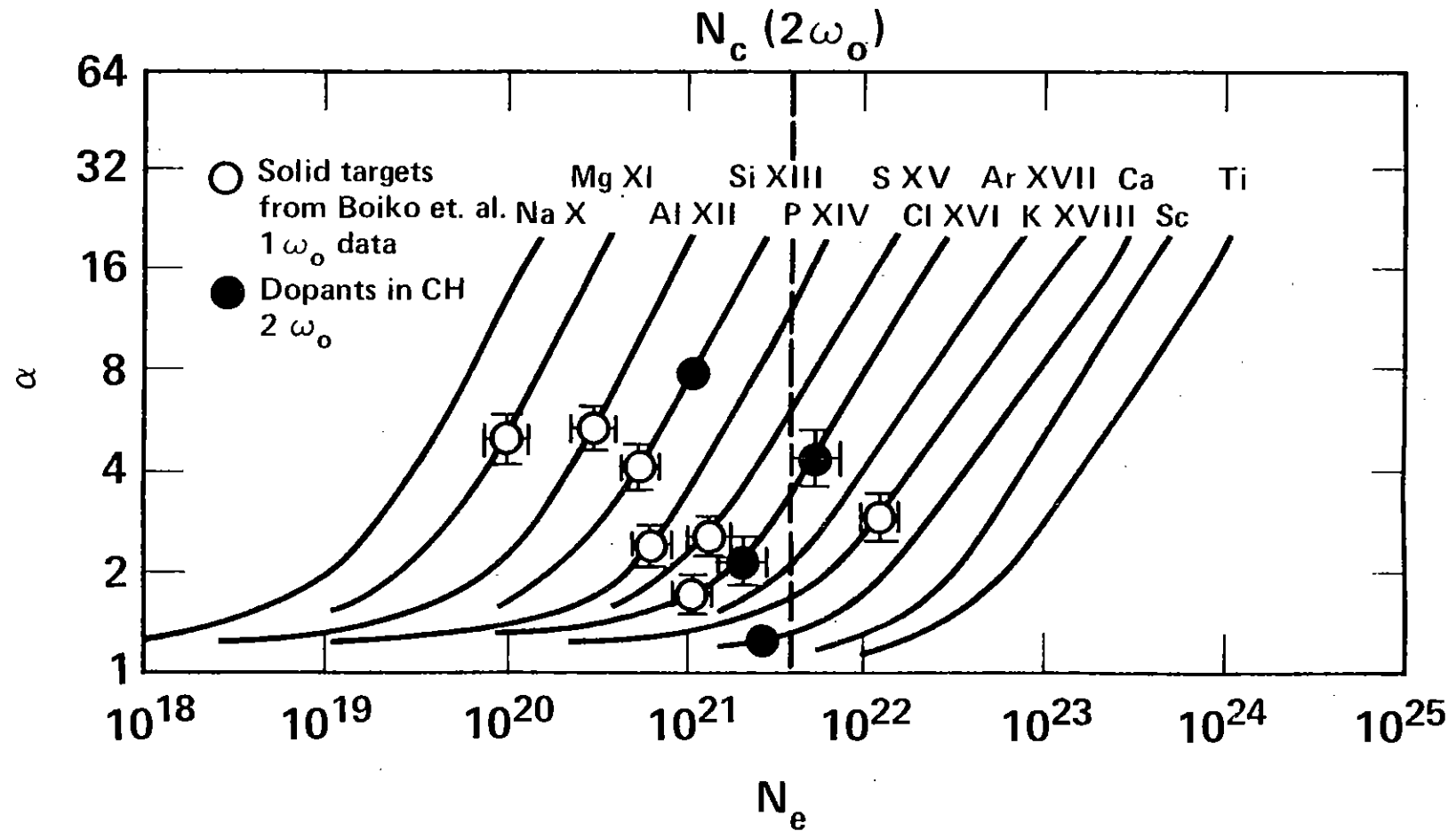


FIG.7.49.
SIMPLE ANALYSIS OF $^1P_1/^3P_1$ RATIO SHOWS THAT DOPANT
EMISSION IN CH OCCURS AROUND CRITICAL DENSITY
FOR 0.532 nm LIGHT



yields the 'radial structure function', which exhibits a series of peaks at $r = r_j + a_j$ where the shift a_j arises because of the omission of the k dependent phase shift $\phi(k)$ in the Fourier transform.

In practice, as the range of coherent electrons in solids is limited to a few Angstroms, the spectrum is usually dominated by contributions from the nearest neighbour atoms, and thus the first scattering shell usually predominates.

There are many advantages to using laser-produced plasmas (LPP) as X-ray sources for EXAFS studies, compared with either conventional X-ray sources, or more recently, synchrotron radiation sources. Firstly the brightness of LPP sources can be several orders of magnitude higher than synchrotron sources. The complete EXAFS spectrum can be recorded in a single 1 ns laser pulse. Secondly it is possible to investigate samples that would be damaged, or decompose, during the repetitive X-ray exposure that is characteristic of a synchrotron source. In the 1 ns necessary for a spectrum to be recorded using a LPP, the sample does not have sufficient time to react. Thirdly, and perhaps most importantly, LPP sources offer the unique possibility of performing time resolved measurements of transient structures, induced by the laser itself, or some other synchronised external stimulus.

In this section we shall report time integrated measurements of aluminium foils, both cold and radiatively heated. Section 7.9.3 will deal with time resolved measurements which utilise the rapid radiative heating of a foil during the 1 ns laser pulse, to show the consequent reduction in amplitude of the characteristic EXAFS modulations.

7.9.2 Experiment

The arrangement used to record static EXAFS spectra is shown in Figure 7.50. A ~ 100 J, 1 ns, 1.06 μm laser pulse was frequency doubled to provide 30-40 J of 0.53 μm light. This was focused down to ~ 200 μm spot sizes onto various target materials for aluminium EXAFS, a suitable candidate for trial experiments. Bismuth was chosen as the best X-ray

source, as the spectrum above the K-edge of Aluminium is almost completely featureless. The dispersing crystal, EXAFS samples, and X-ray film were all housed in a Mk 2 minispectrometer. Several crystals were used, notably ADP, and PET.

The double slit arrangement on the front of the minispectrometer allowed several possibilities for recording multiple spectra on a single laser shot. Initially both the incident and transmitted spectra were recorded, by putting a 2 μm aluminium foil over only one half of the film. This allowed the incident X-ray spectrum to be checked for any undesirable line emission in the region of interest. Subsequently however, double foil shots were recorded involving two Aluminium foils, one over the film plane to provide a 'cold' spectrum, and a second at varying distances > 1 mm from the plasma. This second foil was heated throughout its depth by soft X-rays such that its temperature was rapidly raised during the 1 ns laser pulse. Calculations suggest that the temperature rise is of the order of several hundreds of degrees centigrade, possibly sufficiently high to reach the melting point of Aluminium.

Figure 7.51 shows the EXAFS spectrum of a 2 μm Aluminium foil obtained using Iron and Bismuth targets. It is clear that whereas Iron shows extensive line radiation, Bismuth is free from spectral features in this region. The structure present in the Iron spectrum has been assigned to Fe XXIV emission lines, according to the tabulations by Boiko (7.61). They serve incidentally as a wavelength calibration for the EXAFS data using Bismuth targets, as the spectrometer position was identical on the two consecutive shots in Figure 7.51.

It should be pointed out that the quality of the spectra of Figure 7.51 is as good, if not better than, those obtained from Aluminium using synchrotron sources. Using the ACO storage ring at Orsay, Fontaine et al (7.62) took 1.5 h to record a similar spectrum. The feature marked 'b' however in Figure 7.51 was not as well resolved in their spectrum, and the feature 'a' was not observed at all.

Figure 7.52 shows EXAFS spectra of both 'hot' and 'cold' Aluminium. The

FIG. 7.50 MINISPECTROMETER ARRANGEMENT USED TO RECORD LASER EXAFS SPECTRUM

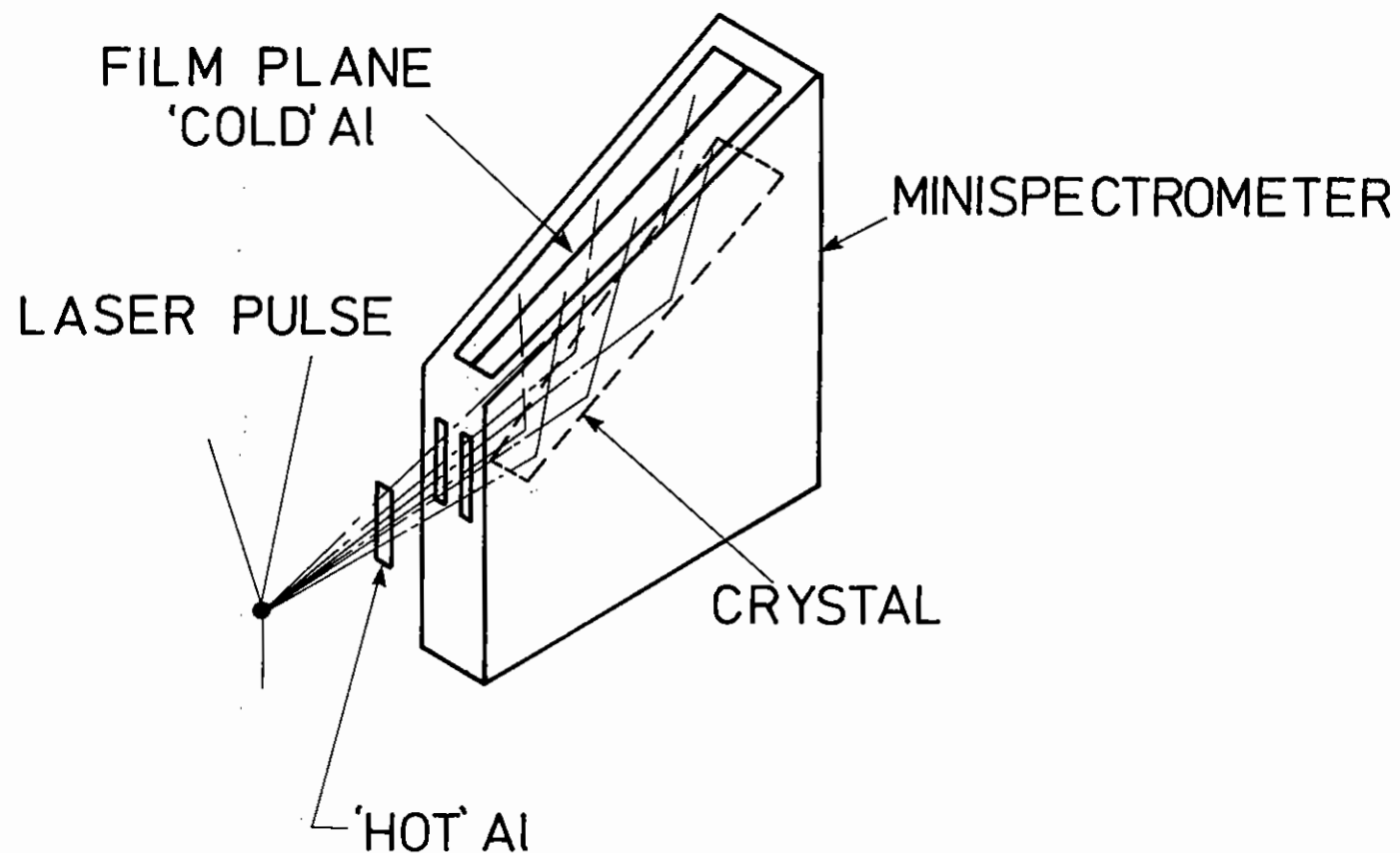


FIG. 7.51. TIME INTEGRATED EXAFS SPECTRA OF 2 μ Al
FOIL WITH DIFFERENT X-RAY BACKLIGHTING
SOURCES

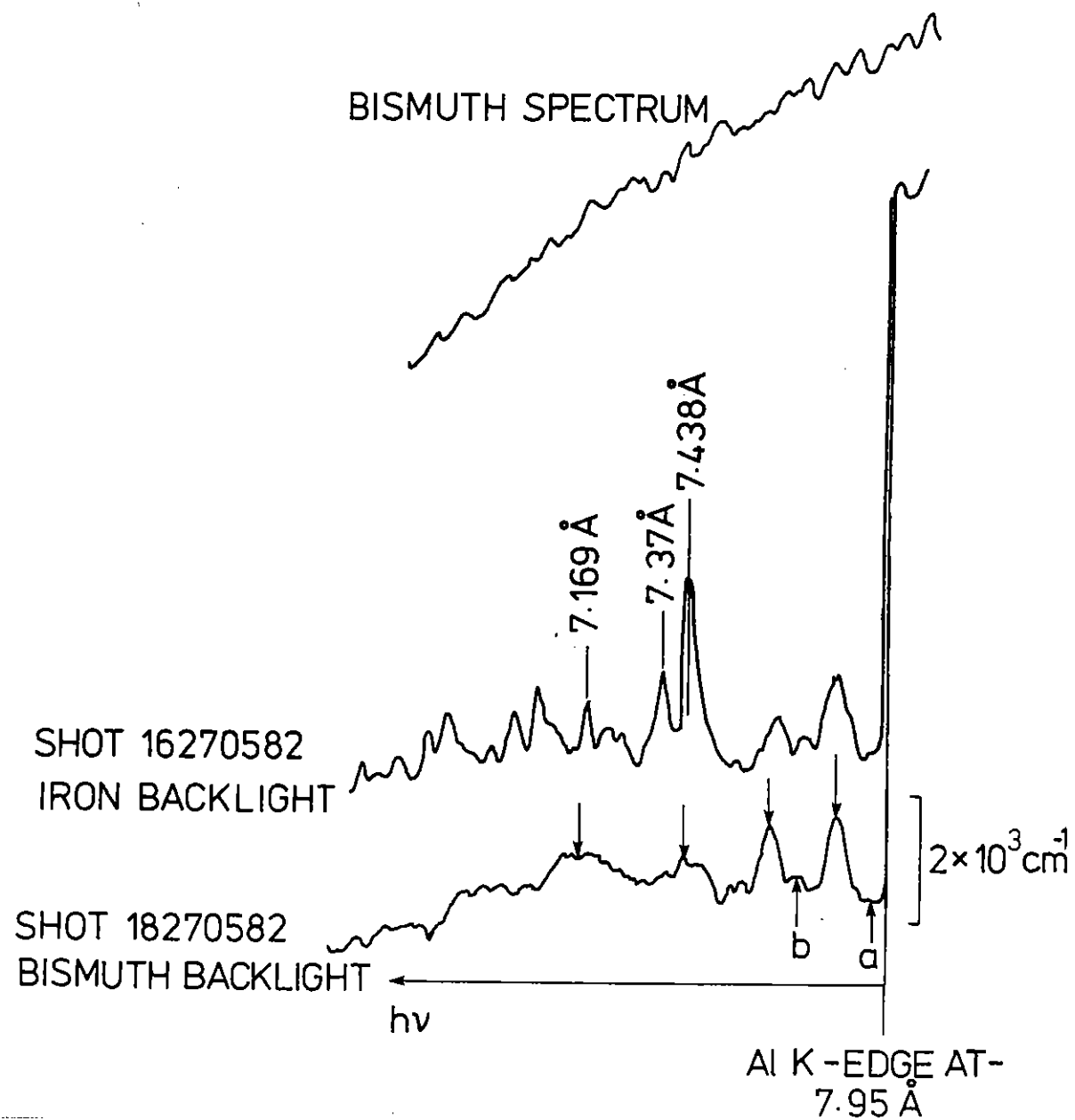
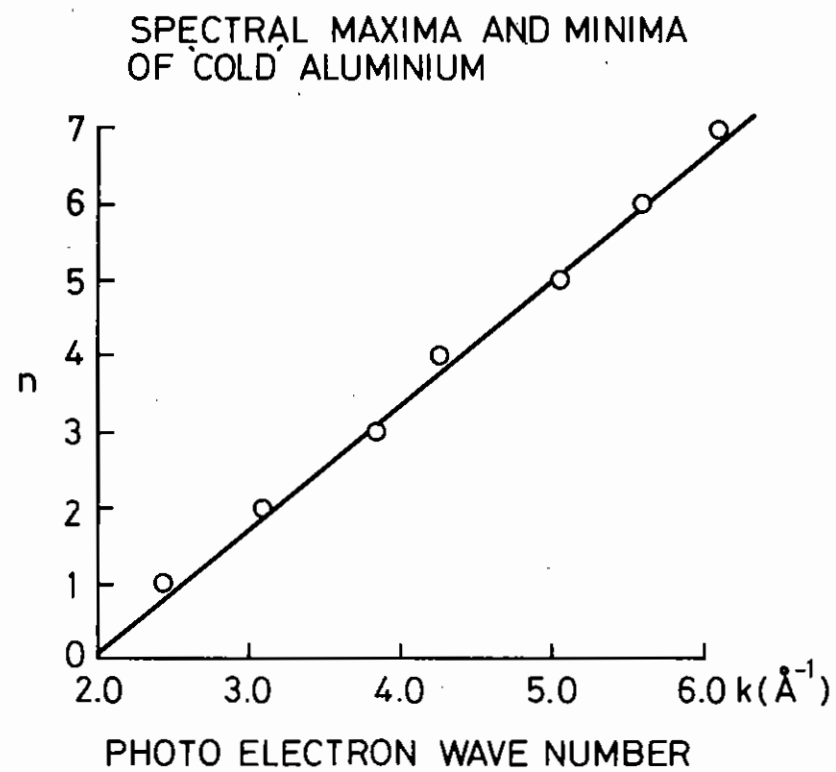
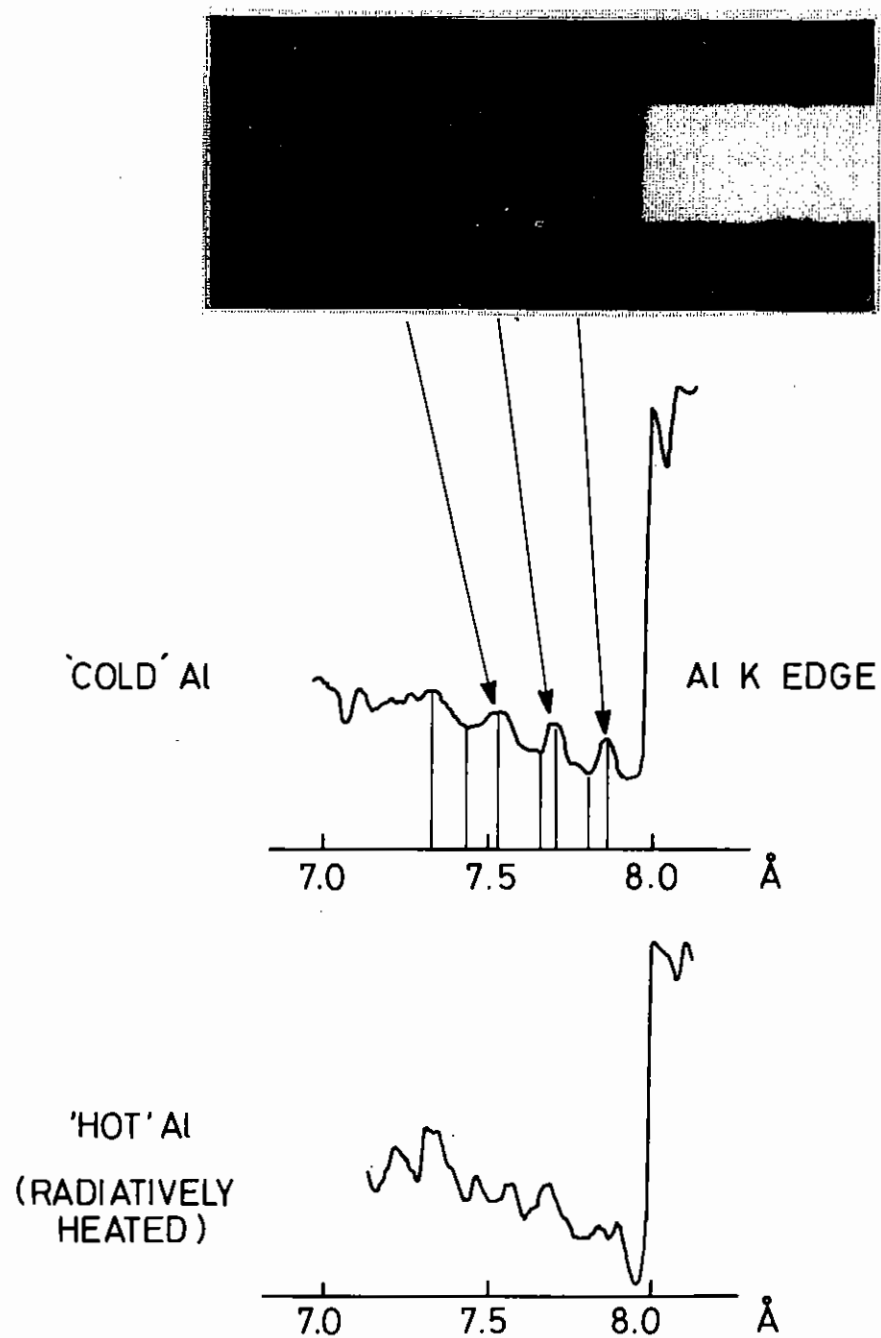


FIG. 7.52 LASER EXAFS SPECTRUM OF ALUMINIUM



ATOMIC SPACING DERIVED FROM EXAFS
 $= (R_1 - \alpha_1) \approx 2.5 \text{\AA}$



fine structure for 'hot' Aluminium is clearly very different to the cold spectrum. This would suggest that the radiative heating has substantially altered the room temperature FCC structure of the Aluminium lattice. Fourier transforms of both spectra would provide information on the time-integrated behaviour of the lattice. This is in progress, using EXAFS codes developed at Daresbury Laboratory (7.63).

In the absence of a Fourier transform, atomic positions can still be determined however, as for certain lattice structures, the EXAFS spectrum is dominated by the nearest neighbour shell. The position of the principal maxima and minima therefore may be plotted sequentially as a function of k , as outlined in (7.64). This is shown in Figure 7.52 for the first seven maxima and minima. The straight line obtained is given by

$$n\pi = 2k(R_1 - \alpha_1) + 2\beta_1$$

from which a value of $R_1 - \alpha_1 \approx 2.5 \text{ \AA}$ is obtained. This is in very good agreement with the value of $2.45 \text{ \AA}$ reported in (7.62) for the peak in the radial structure function obtained from their EXAFS spectrum via a Fourier transform.

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7.9.3 Time resolved EXAFS

A recent preliminary experiment has demonstrated the feasibility of making time resolved EXAFS measurements with a resolution $< 100 \text{ ps}$ using a LPP source and a spectrally dispersive X-ray streak camera.

The experimental arrangement is shown in Figure 7.53. 30-40 J, 1 ns laser pulses of $0.53 \text{ \mu m}$ wavelength were focused down to $\sim 150 \text{ \mu m}$ diameter spots onto a bismuth backlighting target. The plasma thus produced provided both the backlighting spectrum and the heating source for this experiment. An aluminium foil (typically 2-3 $\text{\mu m}$ thick) was placed between 1 and 12 mm from the plasma region, this distance being varied from shot to shot, to set the rate of heating of the foil. The backlighting spectrum viewed

through the foil was recorded on the streak camera via a flat TLAP ($2d = 25.9 \text{ \AA}$) Bragg diffracting crystal.

A typical spectrum recorded through a sample $> 1 \text{ cm}$ distant from the plasma is shown in Figure 7.54 with the unfiltered background spectrum from a different shot for comparison. Modulations in the filtered spectrum to the high energy side of the K edge can be seen, and correspond to the EXAFS bands seen on the time integrated measurements reported in the previous sub-section. Figure 7.54 also shows the spectrum recorded through an absorber situated $\sim 3 \text{ mm}$ from the plasma. Microdensitometer tracings at various times during the pulse are shown alongside. It can be seen that the depth of the EXAFS modulations decreases during the pulse as the foil heats up, as has been seen with the previous time integrated spectra, implying that correlation in the aluminium lattice is being lost on a timescale comparable to the heating time. A qualitative (and predictable) result is that the photoelectrons liberated into the lattice by the X-ray heating radiation equilibrate with the ionic lattice in a time $<$ the laser pulse length. Detailed analysis of these temperature induced changes is to take place, using the EXAFS codes at Daresbury (7.63).

The main result of this experiment, however, is a demonstration that single shot time-resolved EXAFS measurements on a sub nanosecond timescale can be made using a laser plasma backlighting source. Such a source is ideally suited to time resolved studies where the time dependent event can easily be synchronised to, or triggered by, the laser.

D Bradley (Imperial College)

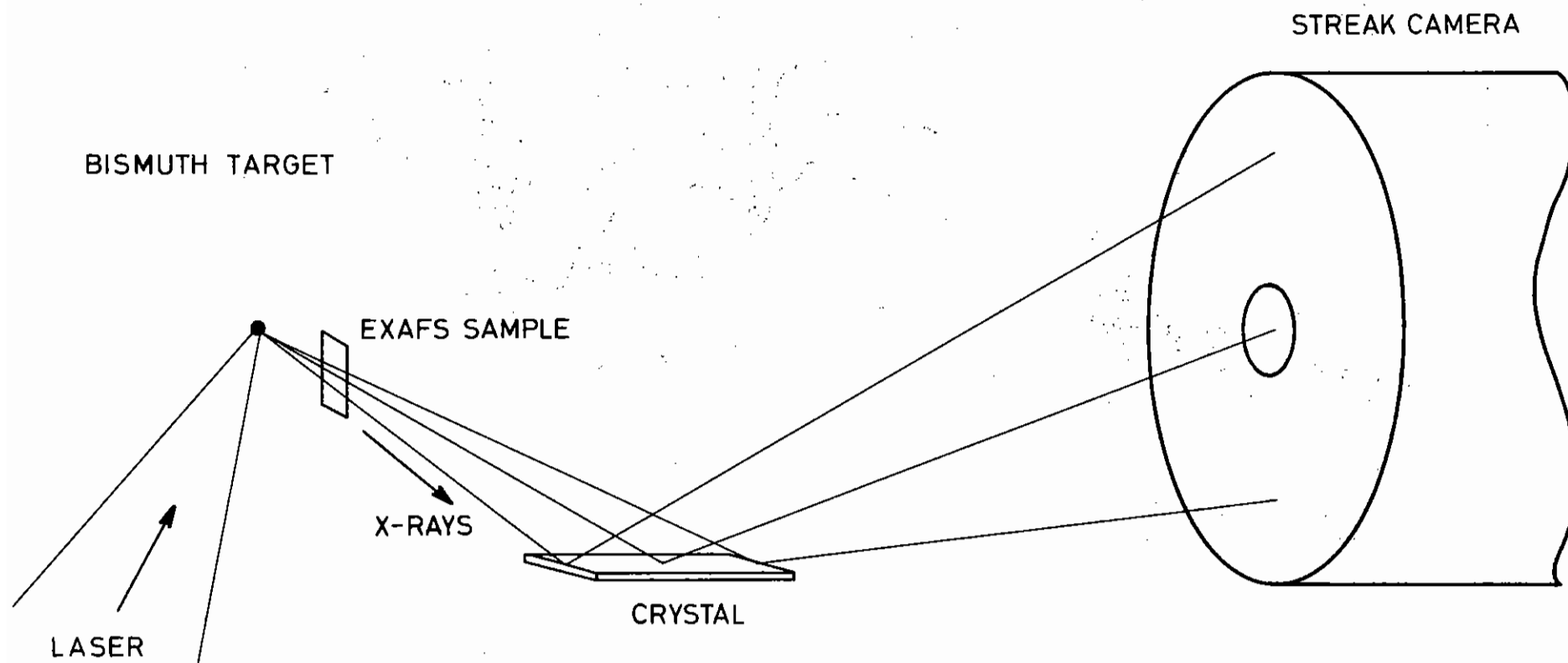


FIG. 7.53. EXPERIMENTAL SET-UP FOR STREAKED EXAFS

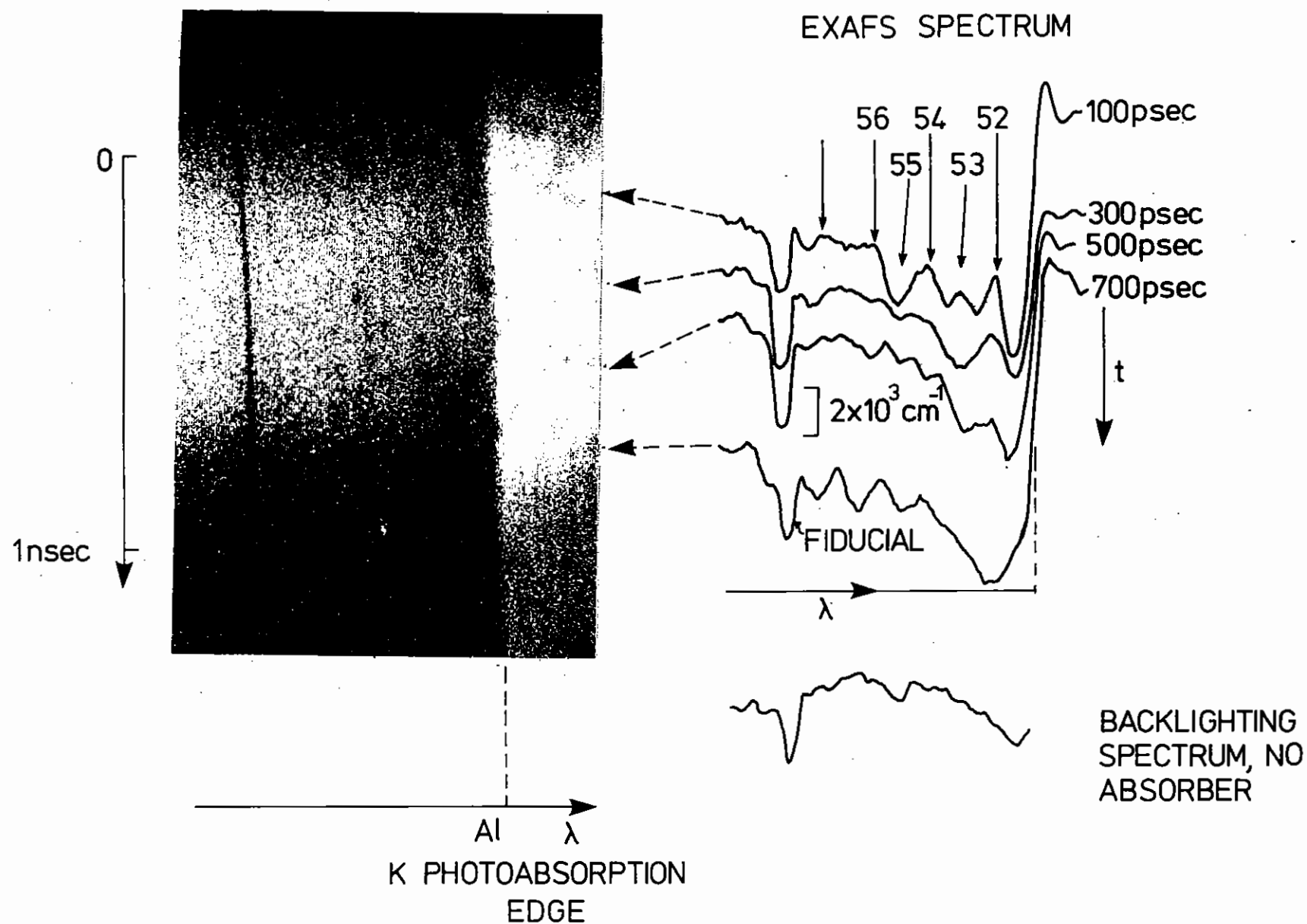


FIG.7.54 TIME RESOLVED EXAFS SPECTRUM OF Al

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CHAPTER 8 THEORY AND COMPUTATION

INDEX

- 8.1 INTRODUCTION page 8.1
- 8.2 COMPRESSION STUDIES page 8.1
 - 8.2.1 Fluid Code Development
 - 8.2.2 Rayleigh Taylor Instability
 - 8.2.3 Effects of Heat Flux Limitation of Thermal Smoothing
 - 8.2.4 An Analytic Model of Ion Beam Heated Targets
 - 8.2.5 Soft X-ray Preheat as a Constraint on Ablative Implosions
- 8.3 LASER PLASMA INTERACTIONS page 8.29
 - 8.3.1 Oscillations Due to Feedback in Interaction Processes
 - 8.3.2 The Use of the Finite Element Method in Electrostatic Particle Codes
 - 8.3.3 The Effect of Magnetic Field on Raman Sidescatter
 - 8.3.4 Kinetic Theory of Raman Scatter in a Magnetized Plasma
 - 8.3.5 The Effect of Brillouin Scatter on Raman Scatter
- 8.4 TRANSPORT PROCESSES page 8.47
 - 8.4.1 Grad's 13 Moment Equations and Electron Thermal Transport
 - 8.4.2 Nonlinear Heat Flow
 - 8.4.3 Time-independent Solution of the Vlasov-Fokker-Planck Equation for a Planar Ablating Plasma
 - 8.4.4 Heat Flow Instabilities
 - 8.4.5 Radiation Transport with a Planckian Distribution
 - 8.4.6 LTE Opacity Calculations
 - 8.4.7 The Calculation of Radiation Preheat in Laser Driven Implosions

8.4 Cont

- 8.4.8 Some Numerical Aspects of NLTE Radiation Transfer Calculations
- 8.4.9 1982 CECAM Workshop on 'The Flux Limiter and Heat Flux Instabilities'
- 8.4.10 The CECAM Workshop on 'Non-LTE Radiative Transfer'

REFERENCES page 8.77

CHAPTER EDITOR: A R Bell

CHAPTER 8 THEORY AND COMPUTATION

8.1 Introduction

Two important growth areas in theory are two-dimensional hydrodynamic effects (both Rayleigh-Taylor instabilities and thermal smoothing) and radiation transport calculations. There are two main objectives of the latter, the calculation of pre-heat and its effect on compression, and as a diagnostic for comparison with emission spectra. As closer comparison with experiment is sought, the coupling between scattering processes or a scattering process and the ponderomotive force steepening is being considered; in the former case frequency shifts and in both cases new instabilities can arise. One dimensional heat flow and the understanding of the flux limiter is now a mature subject, and the effect of ion motions is now being included. The stability of non-linear heat flow to electromagnetic perturbations is now being studied from the Fokker Planck equation, and this theory might account for the experimental filamentary structure.

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8.2 Compression Studies

8.2.1 Fluid Code Development

8.2.1(a) Quasi-Lagrangian Rezoning Techniques

The aim of this programme is the development of accurate, conservative hydrodynamic algorithms capable of following the fluid yet preserving an orthogonal structure. Three essential differences characterise these models from more conventional fully Lagrangian codes, namely the scheme is strongly (exact to round-off) energy conservative, orthogonally grided and does not use an artificial viscosity, each improving the accuracy of the code. In fully Lagrangian codes, the one dimensional von Neumann-Richtmyer algorithm is weakly $[O(Dt^3)]$ conservative, but two dimensional (eg Shultz) only consistently $[O(Dt)]$, this may give rise to a numerical

energy generation which, in extreme cases, is of the same order of magnitude as the laser sources.

The use of an orthogonal grid eliminates the possibility of pathological zones. A well ordered mesh, however, is achieved at the expense of the ability to simply follow contact surfaces, although procedures are available for this purpose (8.01). The necessary mesh/fluid slip (rezoning) allows the possibility of mesh modification to meet other criteria, for example a maximal time-step.

Inviscid hydrodynamics is purely dissipationless. In a real fluid dissipation in the form of viscosity and thermal conduction causes the necessary entropy increase across a shock discontinuity. In the absence of dissipation the fluid behaves reversibly and the shock discontinuity must be followed by oscillations. In most routines, eg Von Neumann-Richtmyer, Shultz, Lax-Wendroff this is accomplished by introducing an artificial viscosity which provides the necessary dissipation over distances of the order of the cell width in a compressing flow. In comparison with real dissipation such effects are clearly much more coarse grained, and in consequence can lead to non-physical effects. As an example we may identify three such in the case of a spherically convergent shock driven by a piston; (i) the entropy jump at the centre is greatly in excess (ii) entropy generation occurs in the adiabatic compression following the incoming shock (iii) the incoming shock moves too rapidly by virtue of lower density and higher temperature between the shock and piston. In two dimensions anomalous energy transport across long thin zones occurs due to artificial viscosity.

The basic approach to this problem using cell centred grids is described in a recent paper (8.02) using a novel stable fluid algorithm and flux corrected transport. In this report we outline more recent work on this topic.

8.2.1(b) Conservative Non-uniform Advection Procedures

The conventional approach to advection considers the flow of uniform fluid

on a uniform mesh. We have shown that a general second order accurate conservative finite difference representation on a non-uniform mesh of the equation:

$$\frac{\partial \rho}{\partial t} = - \frac{\partial}{\partial z} (\rho u) \quad (1)$$

may be written

$$\rho_i \Delta z_i = \rho_i^0 \Delta z_i^0 - q_{i+\frac{1}{2}} + q_{i-\frac{1}{2}} \quad (2)$$

where ρ_i (ρ_i^0) and Δz_i (Δz_i^0) are the final (initial) density and cell width. The flux across the face $i + \frac{1}{2}$

$$q_{i+\frac{1}{2}} = \frac{1}{2} Q_{i+\frac{1}{2}} \{ (\rho_{i+1}^0 - \rho_i^0) - Q_{i+\frac{1}{2}} (\rho_{i+1}^0 - \rho_i^0) \} \cdot \Delta z_{i+\frac{1}{2}}^{10} \quad (3)$$

where $\Delta z_{i+\frac{1}{2}}^{10}$ is the initial mesh spacing ($z_{i+1}^0 - z_i^0$),

$$Q_{i+\frac{1}{2}} = v_{i+\frac{1}{2}} Dt / \{ \Delta z_{i+\frac{1}{2}}^{10} + (u_{i+1} - u_i) Dt \} \quad (4)$$

where $v_{i+\frac{1}{2}}$ is the flow velocity of fluid through the mesh at the face $(i+\frac{1}{2})$.

In general any conservative three point advection scheme may be written

$$q_{i+\frac{1}{2}} = \eta_{i+\frac{1}{2}}^- \rho_i^0 - \eta_{i+\frac{1}{2}}^+ \rho_{i+1}^0 \quad (5)$$

In flux corrected transport (FCT) an additional diffusion is added to the transport stage (8.03) which is subsequently removed by extremum limited anti-diffusion. In order to match equation (2) by the complete routine

$$q^- = \frac{1}{2} Q(1 + Q) + \eta : q^+ = -\frac{1}{2} Q(1 - Q) + \eta \quad (6)$$

where η is the prescribed diffusion coefficient. Thus we have

Shasta

$$\eta = \frac{1}{8} \quad q^- = \frac{1}{2}(\frac{1}{2} + Q)^2 \quad q^+ = \frac{1}{2}(\frac{1}{2} - Q)^2 \quad (7)$$

Donor cell

$$\eta = \frac{1}{2}[Q](1 - [Q]) \quad q^- = \text{Max}(\nabla, Q) \quad q^+ = \text{Min}(\nabla, Q) \quad (8)$$

The advection may be calculated by one of three methods (8.03), thus if ρ^D is the density after diffusion and ρ^T due to transport alone

$$\rho_i^T = \rho_i^D - \eta_{i+\frac{1}{2}}(\rho_{i+1}^D - \rho_i^D) + \eta_{i-\frac{1}{2}}(\rho_i^D - \rho_{i-1}^D) \quad (9)$$

we have the density after advection

$$\begin{aligned} \rho_i^D - \eta_{i+\frac{1}{2}}(\rho_{i+1}^D - \rho_i^D) + \eta_{i-\frac{1}{2}}(\rho_i^D - \rho_{i-1}^D) & \text{ explicit} \\ \rho_i^A = \rho_i^D - \eta_{i+\frac{1}{2}}(\rho_{i+1}^T - \rho_i^T) + \eta_{i-\frac{1}{2}}(\rho_i^T - \rho_{i-1}^T) & \text{ phoenetical} \\ \rho_i^D - \eta_{i+\frac{1}{2}}(\rho_{i+1}^A - \rho_i^A) + \eta_{i-\frac{1}{2}}(\rho_i^A - \rho_{i-1}^A) & \text{ implicit} \end{aligned} \quad (10)$$

Of these implicit does not treat shocks well and will not be considered further.

In a Lagrangian system no fluid flows through the mesh, $Q = 0$. Thus there is a residual term for explicit FCT, but only clipping due to the extremum limiting for phoenetical. With explicit FCT this is a high order dissipation for on a uniform system the mode e^{ikz} has an amplification factor

$$1 - \eta^2[1 - \cos(k\Delta z)]^2 \quad (11)$$

ie a fourth order damping. This term was analysed in detail for the case $\eta = \frac{1}{8}$ in our earlier work (8.02). The existence of this damping term as an alternative to artificial viscosity for providing the required dissipation even in purely Lagrangian systems was proposed and shown to be satisfactory (8.02) With phoenetical anti-diffusion some fourth order damping will still occur as a consequence of the clipping.

8.2.1(c) Flux-correction as an artificial dissipation

The fourth order residual damping considered as an alternative to artificial viscosity has a number of features (i) as a fourth order term (rather than second order) it is dominantly short range; (ii) over distances of the order of a cell width a discontinuity has a stable steady structure; (iii) slow compression can occur nearly adiabatically. It was shown in our earlier work that the dissipation from explicit SHASTA was sufficiently strong to eliminate overshoot in the case of a planar one dimensional shock in a Lagrangian system. Subsequent studies have confirmed this result more generally. Modelling of collapsing spherical shocks has shown significant improvement over the von-Neumann-Richtmyer scheme in the accuracy with which the analytic solution can be reproduced. In particular the strong dissipation at the centre, and consequent low density are eliminated, and the adiabatic compression between shock and piston correctly represented.

In principle phoenetical flux correction may also provide dissipation to ensure a satisfactory shock behaviour, however the primary decomposing of the post-shock oscillation is a consequence of clipping similar in concept (though not in method) to van Leer's monotonic approach (8.04). In consequence the shock fronts are more diffuse, in the absence of short range steepening, a result confirmed in numerical tests. In spherical shock tests phoenetical SHASTA performed markedly less well than the explicit algorithm; due essentially to the inability of the clipping to operate satisfactorily in the region of increasing density behind the shock fronts. The performance of this version was comparable with the von Neumann-Richtmyer scheme.

The excellent performance of explicit FCT in shock wave tests is, as is usual, only obtained at a price. We have already remarked that the fourth order dissipation smooths long wavelength disturbances, but actually sharpens very short perturbations: a result which may cause severe error in experiments, or where a strong gradient exists. Thus consider a density profile

$$\rho_1 = \rho_0 x^1 \quad (12)$$

on a uniform mesh. After diffusion/anti-diffusion with a zero velocity ($Q = 0$)

$$\rho_1 = \rho_0 [1 - \text{Min}(\eta, x, 1/x)(x-1)^2/x][1 + \eta(x-1)^2/x]x^1 \quad (13)$$

and the profile is markedly steeper. This problem which is associated with explicit anti-diffusion can be resolved in two ways:

- (a) Phoenetical anti-diffusion.
- (b) Donor-cell FCT, where $\eta = 0$ if $Q = 0$.

Although it appears probable that phoenetical FCT was devised to eliminate this problem (8.03), we have preferred to use donor-cell explicit FCT in cases of expanding or stationary flow where dissipation is not required. Our reasons for this are due to the excellent performance of this algorithm in tests against analytic results for expanding flow (8.02), and the fact that the term minimises the dissipation (and hence irreversibility) for quasi-Lagrangian flow where $Q \approx 0$. Explicit SHASTA is used for compression. Although it can be shown the phoenetical anti-diffusion is well behaved if $\eta < \frac{1}{2}$ comparative tests of an explicit/phoenetical algorithm have not yet been carried out.

8.2.1(d) Rezoning with Fixed Boundaries

In many cases it is required to limit the extent of the mesh, for example by the focal spot size. In this case it is necessary to rezone the mesh back to the boundaries. A simple way of doing this, consistent with cell centred grid, is to add to the grid point velocity at z_1 an extra component $\{(V_b - V_{\frac{1}{2}})(z_f - z_1) + (V_f - V_{N+\frac{1}{2}})(z_1 - z_b)\}/(z_f - z_b)$ where V_b and V_g are the allowed boundary velocities, and $V_{\frac{1}{2}}$ and $V_{N+\frac{1}{2}}$ those deduced from the grid velocity. Such a prescription is not satisfactory for it may lead to continuous rezoning with some cells being reduced to zero

size. To demonstrate this effect and illustrate its solution consider the case where $z_b = V_b = 0$, and suppose, as in a Lagrangian scheme V_1 is a function of z : $V_1 = V(z_1)$. The total grid velocity is then

$$V_1^1 = V_1 + (V_f - V_{N+1/2}) z_1/z_f$$

$$V_1^1 = V(z_1) + [V_f - V(z_f)] z_1/z_f$$

In the case (for example) where V_f is required to be zero, the grid is only stationary if $V(z) = V(z_f) z/z_f$, and furthermore if $\lim_{z \rightarrow 0} [V(z) - V(z_f) z/z_f] < 0$ the motion of the cell centre near $z = 0$ will continuously approach the origin (or if $\lim_{z \rightarrow z_f} [V(z) - V(z_f) z/z_f] > 0$ approach $z = z_f$). In particular the mesh spacing will progressively reduce to zero.

The solution to this problem is to add a restoring velocity to coerce the mesh into a well ordered state, and generate a steady condition. As an example let us require the mesh basis to be a uniform grid where $z_1 = x_1 = (1-1/2) z_f/N^2$, N being the total number of cells. Clearly $z_{N+1/2} = x_{N+1/2} = z_f$. The grid velocity:

$$V_1^1 = V_1 + (V_f - V_{N+1/2}) \{ z_1/z_f - \alpha(x_1 - z_1)/z_f (dz/dx)_1 \}$$

has the desired properties provided

$$\alpha > 1 - \frac{z_f}{[V_f - V(z_f)]} \left. \frac{dv}{dz} \right|_s$$

where z_s is a singular point at which $V(z) = [V_f - V(z_f)] z/z_f$: the most important of these occurring at $z = 0$ and $z = z_f$.

The generalisation to a mesh expanding at both ends is straightforward for if $V_1 < 0$ and $V_{N+1/2} > 0$ there must exist some points at which $V(z) = 0$, and about which each side may, in principle, be treated as above.

Extensive experience in numerical calculations has shown this algorithm to

work extremely well. Generalisations to use different generating distributions by replacing $(x_1 - z_1)/z_f$ by an appropriate alternative are allowed, and indeed the scheme is extremely flexible.

8.2.1(e) Monitoring a Well Structured Mesh

In one dimension the quasi-Lagrangian nature of the rezoning maintains a well ordered mesh by virtue of mass conservation above, particularly if the grid-slip constraint described earlier (8.02) is used. However in two dimensions an incoming flow along one co-ordinate direction may give rise to an outflow in the other eg two jets. In consequence cells in the first direction continuously compress, whilst those in the other expand. Such a problem involves a rotation of the flow and is clearly difficult to handle using Lagrangian dynamics (impossible with a normal two dimensional Lagrangian code). In such circumstances it is important to prevent cells near the origin becoming vanishingly small. As in the previous section an ordering velocity can be used to avoid this effect. As above this can be chosen extremely generally, but will here be based on a uniform mesh. An appropriate correction is to add to the grid velocity, V_1 , an additional component reducing the motion, and coercing it to a uniform grid namely: a net velocity:

$$V_1 = \beta V_1 + (1-\beta) \text{Min}[0, (x_1 - z_1)/(z_f - z_b)] \text{Sign}(V_1)/Dt$$

where an appropriate value of the constant has been found to be:

$$\beta = \text{Min}[1, 10 dz/dx]_1$$

which maintains a reasonable mesh without modifying its motion under normal conditions.

These algorithms described in this report have been included in the code POLLUX. Over the past six months considerable experience of their behaviour has been obtained in modelling carbon fibres heated by one and four beam laser irradiation. The code has proved extremely flexible coping with both consequent collapse with shearing in the four beam flow

and shock propagation through the fibre in one.

8.2.1(f) Hybrid Code Development

The hybrid code is a simple hydrodynamic model for studying three dimensional flows using essentially one dimensional techniques. The model is based on an exact solution to the equations of hydrodynamics in terms of similarity solutions in two dimensions.

The range of the model is greatly extended by the use of a local approximation whose accuracy can be monitored, to describe a wide variety of expanding flows of heated cylindrical bodies. Although the relevant equations can be directly integrated by the method of characteristics, this approach is complicated and difficult to apply. For this reason a simpler method based on the von Neumann-Richtmyer Lagrangian algorithm has been adopted.

Over the past year the accuracy of the representations generated by the model has been studied in some detail. Since the model uses the similarity solutions, whose behaviour was examined in earlier work (8.05) it is to be expected that a similar level of accuracy would be achieved in the present case. The tests were carried out by comparison with runs of the two dimensional code POLLUX with an isothermal structure closely reproducing that required by the model. Since the hybrid model is used to represent expansions only, explicit donor cell FCT was used (SHASTA giving obvious errors), and provided a useful test of the quasi-Lagrangian rezoning routines.

One of the few closed solutions of the hybrid equations enables one to calculate the propagation of the first characteristic along a rod of pre-heated gas as it expands laterally. The solution is of some intrinsic interest as it demonstrates that the characteristics will always penetrate through the rod to its centre, and that the entire fluid will achieve some lateral motion, except at the plane of symmetry. The actual trajectory of the characteristic is strongly dependent on the value of the similarity parameter, Γ . This calculation thereby provides a sensitive test of the

accuracy of this parameter by comparison with the two dimensional code runs. It was found that the optimum value of $\Gamma = 0.514$ was in good agreement with our earlier work (8.05).

Compressions of the detailed structure of the density and velocity in expansion of preheated cylinders of aspect ratio greater than about 3 confirm the applicability of the model for these expansions. It is found that the model yields a progressively poorer representation as the aspect ratio approaches unity, but the agreement is, even then, no more than the corresponding sphere/intensity model picture. Further details of the hybrid model and its accuracy will be given in a forthcoming publication.

We may qualitatively identify the optimum value of Γ by noting that initial block profile (uniform density) has a different entropy from that of the final similarity expansion. For the conditions of these calculations (cylindrical expansions) the entropy difference between self-similar and block profiles is:

$$MR_g \{1 + z \ln \Gamma\}$$

initially and increases to MR_g at the end of the block expansion phase, where M is the gas mass and R_g the gas constant. Since the motion, although isothermal, is essentially isentropic, the similarity flow will only be on the correct adiabat if Γ is chosen to match to the entropy of the block, ie if Γ is chosen so that average entropy difference over the block flow is zero. Since the initial difference is zero if $\Gamma = 0.606$, the empirically determined value $\Gamma = 0.514$ must approximately match this condition.

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8.2.2 Rayleigh Taylor Instability

The 1982 annual report (8.06) contained a description of a series of numerical simulations of the Rayleigh Taylor instability in laser accelerated targets and only a small amount of additional information is

described here.

The Rayleigh Taylor instability is an exponential growth from some small "seed" and this may take the form of a perturbation of the target structure or of the laser beam irradiation. If the latter method is chosen the true exponential growth is modified by a secular term representing the spatial variation in ablation pressure. In practice the secular term can be very small since the non-uniform illumination may be smoothed out in the low density ablation plasma. This smoothing is a time dependent process and early in the laser pulse very little smoothing exists. The non-uniform ablation pressure gives rise to a non-uniform shock velocity through the target and vorticity is introduced into the fluid flow. The vorticity grows in time due to the Rayleigh Taylor instability even though the ablation pressure is subsequently smoothed out. This imprinting of the non-uniform illumination depends on the rate at which the smoothing occurs and will be a function of laser irradiance and wavelength as well as of the transverse perturbation scale length.

A few simulations have been made with the POLLUX computer mode (8.07) (8.08) of flat targets accelerated by lasers of 0.53 μm and 0.265 μm wavelength. The target dimensions and transverse perturbation wavelength were held constant. Figure 8.01 shows, for the four simulations, the growth in time of the ripple amplitude in the target. At a given time the simulations at $(2 \times 10^{13} \text{ W cm}^{-2}, 0.53 \mu\text{m})$ and $(10^{14} \text{ W cm}^{-2}, 0.265 \mu\text{m})$ lie close together, as do the simulations at $(10^{14} \text{ W cm}^{-2}, 0.53 \mu\text{m})$ and $(5 \times 10^{14} \text{ W cm}^{-2}, 0.265 \mu\text{m})$. The lower irradiance and shorter wavelength both act to increase the effective seed amplitude of the Rayleigh Taylor instability. Since a factor of two decrease in wavelength is in each case recovered by a factor of five increase in irradiance we can say that this imprinting effect scales roughly as $I\lambda^x$ where $x = \frac{\ln 5}{\ln 2} = 2.3$, at least over the parameter range which has been simulated. This is very much in agreement with simulations in 1D where the rate of expansion of the critical density surface also scales roughly as $I\lambda^2$ and it is this

separation which largely controls the rate of transverse smoothing (8.09).

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8.2.3 Effects of Heat Flux Limitation on Thermal Smoothing

The symmetry which can be obtained in a spherical implosion depends both on the symmetry with which the target is illuminated and the degree to which the plasma "smooths out" the non-uniform illumination. The first set of general calculations on thermal smoothing within the plasma atmosphere were by Gardner and Bodner (8.10) who combined a hydrodynamic model for the radial plasma motion with a simple analytic model for temperature diffusion. In the absence of sources and sinks of energy, and neglecting the $T^{5/2}$ dependence of the plasma thermal conductivity, the temperature T obeys the Laplace equation

$$\frac{\partial^2 T}{\partial x^2} + \frac{\partial^2 T}{\partial y^2} = 0 \quad (1)$$

If x is the axial coordinate and y the transverse coordinate then a perturbation with the form $T = T(x) e^{iky}$ must have an axial dependence $T(x) = T_0 e^{-kx}$, which is the smoothing function used by Gardner and Bodner. The Laplace equation for the perturbed temperature T is a great over simplification and several modifications may be made. For instance the ablation flow may be modelled by using as the unperturbed equation

$$\text{div } \underline{Q} = -V C_V \frac{\partial T}{\partial x}$$

where V is the flow velocity and C_V the specific heat of the plasma. Perturbing this equation gives the dispersion relation

$$k_x^2 + \frac{k_y^2}{L} - k_y^2 = 0(2), \text{ where } L = \frac{\kappa}{C_V V}$$

is a length scale associated with the flow. For large values of k_y this

RAYLEIGH TAYLOR INSTABILITY - MODULATED BEAMS

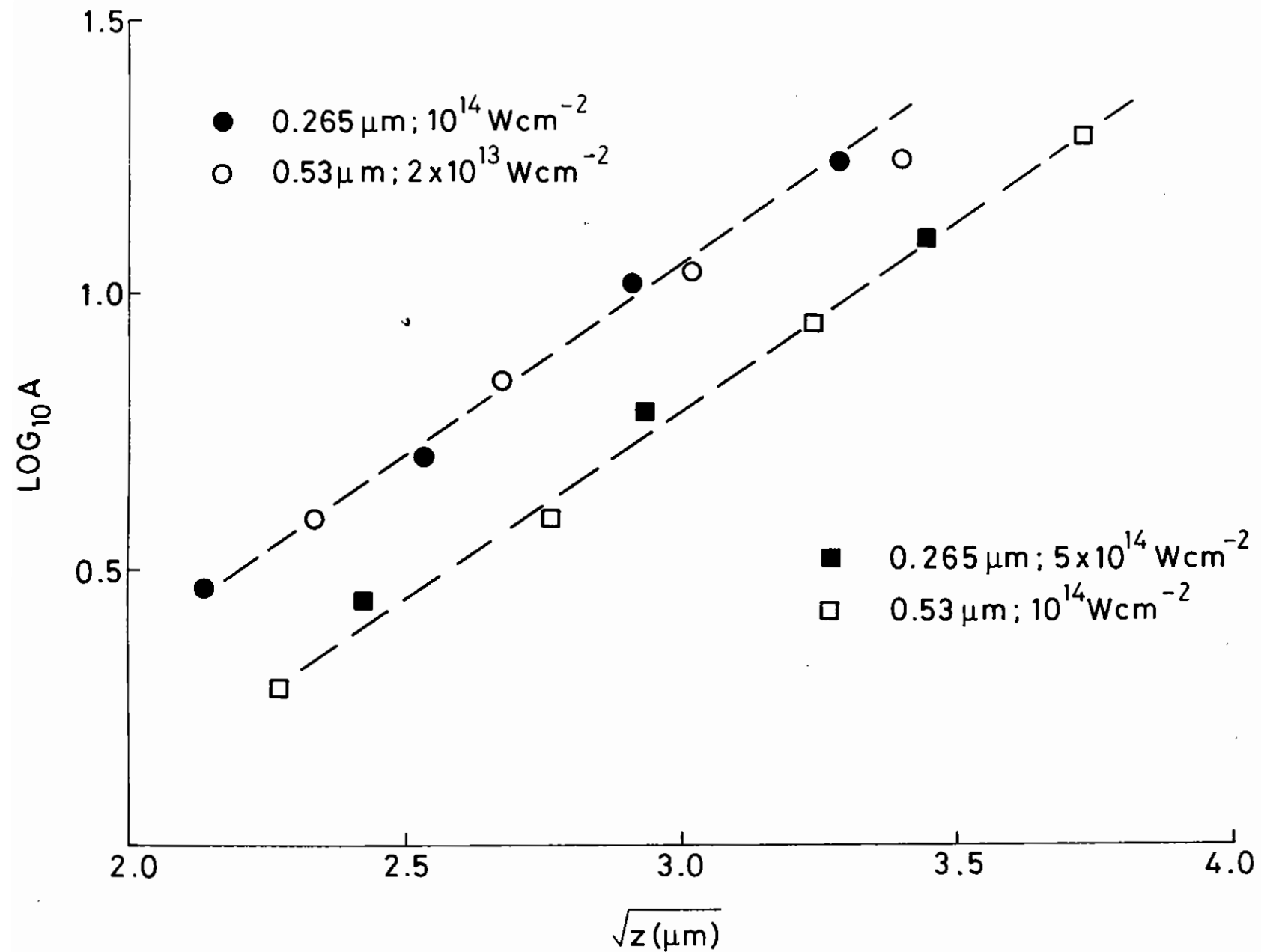


Fig. 8.01 Plot of perturbation amplitude against distance travelled.

gives the previous result while for $k_y L \ll 1$, $k_x \approx 2/L$.

A more interesting variation is to attempt to include the effects of saturated heat flow and a limiting case of this has been discussed by Skupsky (8.11). Although the understanding of saturated heat flow is limited (8.12) it is modelled by the modified heat flow equation

$$\underline{Q} = \frac{\kappa \nabla T}{1 + \frac{|\kappa \nabla T|}{Q_s}} \quad (3)$$

where Q_s is the "saturated" heat flow. In order to analyse this case properly we must include the $T^{5/2}$ dependence of κ and the $T^{3/2}$ dependence of Q_s . Taking the unperturbed equation to be $\text{div } \underline{Q} = 0$ (ie neglecting the ablation flow) the dispersion relation is

$$-k_y^2 + k_x^2 \left(1 + \frac{Q}{Q_s}\right) - \frac{k_x}{L} \left(1 + \frac{Q}{Q_s}\right) + \frac{k_x}{L} \left(\frac{5}{2} - \frac{1}{2} \frac{Q}{Q_s}\right) - \frac{1}{L^2} \left(\frac{5}{2} - \frac{5}{2} \frac{Q}{Q_s}\right) = 0 \quad (4)$$

$$\text{where } \frac{1}{L} = \frac{1}{T} \frac{\partial T}{\partial x}, \quad \frac{1}{L^2} = \frac{(\partial^2 T / \partial x^2)}{(\partial T / \partial x)^2}$$

and the saturated limit corresponds to $Q/Q_s = -1$.

Two limiting forms of this are

$$(a) \quad Q \ll Q_s, \quad T^{5/2} \partial T / \partial x = \text{const}$$

$$\text{then } k_x^2 = k_y^2 + \frac{5}{2L^2} \quad (5)$$

which is the modification of Gardner and Bodner's result for a $T^{5/2}$ conductivity.

$$(b) \quad Q = -Q_s: \quad k_x = \frac{L k_y^2}{3} \quad (6)$$

which is the result obtained by Skupsky.

The interesting feature of Skupsky's result is that for $k_y L \gg 1$ the axial attenuation rate is much faster than would be obtained for classical heat flow ie there is more thermal smoothing. Very crudely it could be said that since the axial heat flow is inhibited the lateral heat flow is relatively enhanced. However the separation D from critical density surface to ablation surface decreases as the heat flow becomes inhibited. Even though k_x becomes greater for inhibited heat flow the overall attenuation factor $F = \int_0^D k_x dx$ may be reduced.

These considerations of smoothing effects in the plasma atmosphere have addressed only the thermal mode of energy propagation whereas the true laser plasma problem is also affected by hydrodynamic effects. The hydrodynamic contribution to "smoothing" has been calculated by Manheimer, Colombant and Gardner (8.13) together with the classical $T^{5/2}$ heat flow. Their results show that the "hydrodynamic" smoothing can easily exceed the "thermal" smoothing. In their analysis the perturbed equations are forced to be steady state so that only x and y spatial derivatives remain. This makes the "hydrodynamic" smoothing a pseudo sound wave which satisfies the local dispersion relation

$$c_s^2 (k_x^2 + k_y^2) = \omega^2, \quad \text{with } 2\omega = 0$$

$$\text{ie } k_x = ik_y \quad (7)$$

[Manheimer, Colombant and Gardner neglect the acceleration of the system which would lead also to gravity waves. Some of the gravity modes are however unstable, being the Rayleigh Taylor modes of the system and would not contribute to "smoothing".]

All of the above effects are naturally included in a hydrodynamic simulation, but are not immediately obvious without the insight given by the analytic theory. A few simulations have been performed with the

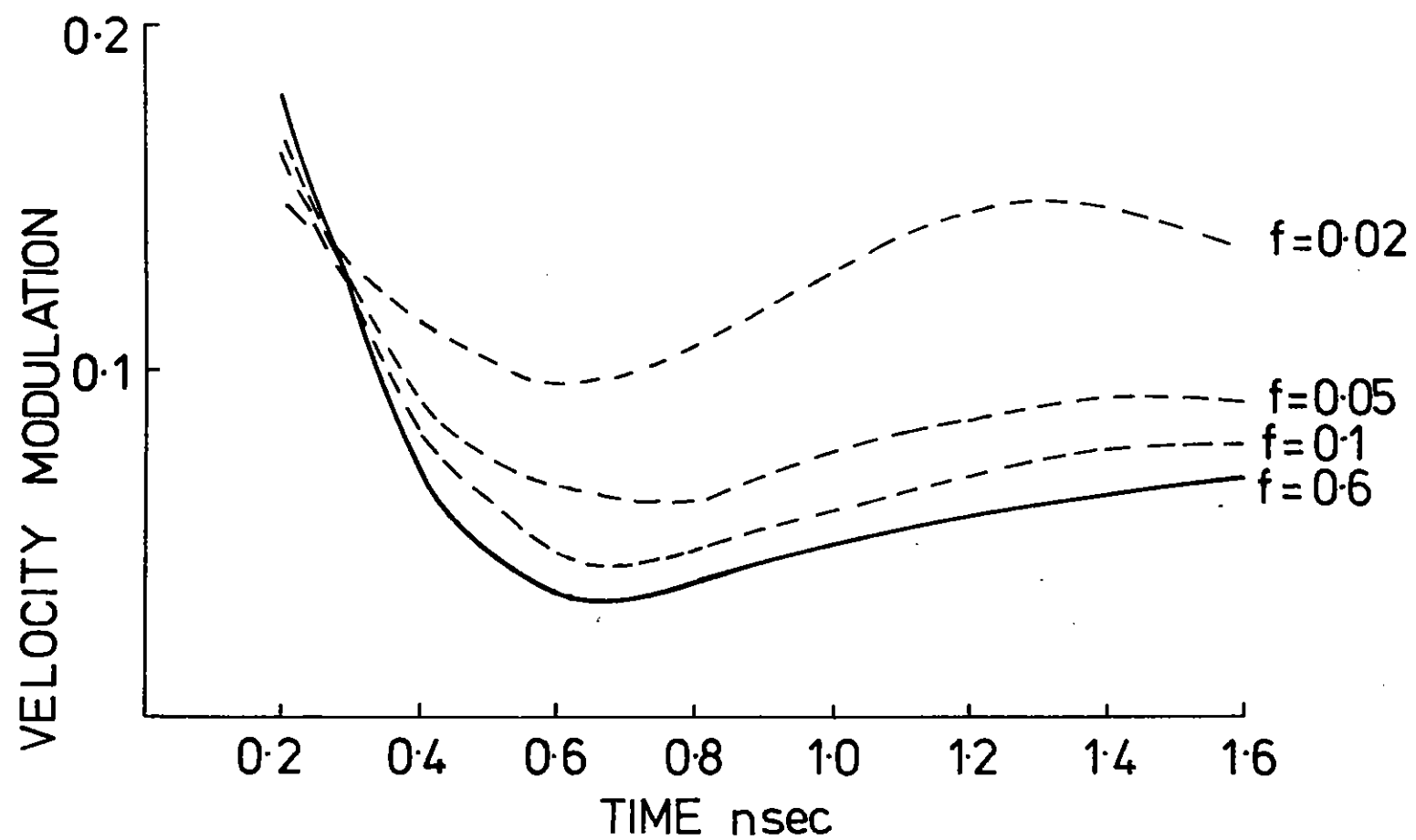


Fig. 8.02 Velocity modulation as a function of time.

POLLUX code (8.07) on the GRAY-1 in order to demonstrate the validity of the theories. A plastic foil 2.5 μm thick is irradiated by a 0.53 μm laser beam at $2 \times 10^{13} \text{ W cm}^{-2}$. The laser is modulated $\pm 20\%$ in intensity with a transverse periodicity $\lambda = 25 \mu\text{m}$. The flux limit in the code is set to various values between 0.02 and 0.6. The modulation of the foil acceleration is given in terms of an average velocity measured on each of the axial mesh lines. Figure 8.02 shows the velocity modulation as a function of time for each of the simulations. Two features are immediately apparent, firstly the smaller values of the flux limit always give a higher modulation of the foil velocity, secondly the velocity modulation shows a tendency to oscillate in time. The oscillation is due to the transverse sound waves excited by the non-uniform illumination. The zero frequency pseudo sound waves of Manheimer et al are a very special case and would only occur in practice if the system could be initialised into an exact eigen function of the coupled thermodynamic-hydrodynamic equations.

In conclusion the problem of symmetry of laser driven implosions is much more complex than might be supposed. To begin with the heat conduction in a laser driven ablation front is not properly understood, secondly Manheimer et al have shown that thermal transport alone does not provide an adequate description of the lateral energy transport and pseudo acoustic modes are important. This work shows that the real problem is also time dependent and should include the effects of acceleration, ie the gravity wave and Rayleigh-Taylor modes. It is probably fair to say that no adequate analytic treatment exists at present and even the numerical simulations have the problem of defining a suitable heat flux vector. For small departures from spherical symmetry the thermally induced magnetic fields are probably unimportant but grossly non-uniform illumination will result in significant modification of the thermal conductivity of the plasma corona.

R G Evans (RAL)

8.2.4 An Analytic Model of Ion Beam Heated Targets

The small spherical pellets proposed for inertial confinement fusion (ICF) applications may be imploded by laser beams, charged particle beams or any other mechanism capable of providing an energy flux of about $10^{14} \text{ W cm}^{-2}$. For a given power density $I \text{ W cm}^{-2}$ the effective pressure generated, p_A and hence the hydrodynamic efficiency of acceleration will depend on the details of the energy deposition from the particular driver. It has become customary to analyse the coupled equations for energy deposition, transport and hydrodynamics using large computer models. While these models give an accurate solution of the particular equations used to describe the system, they do not often provide a good insight into its essential behaviour.

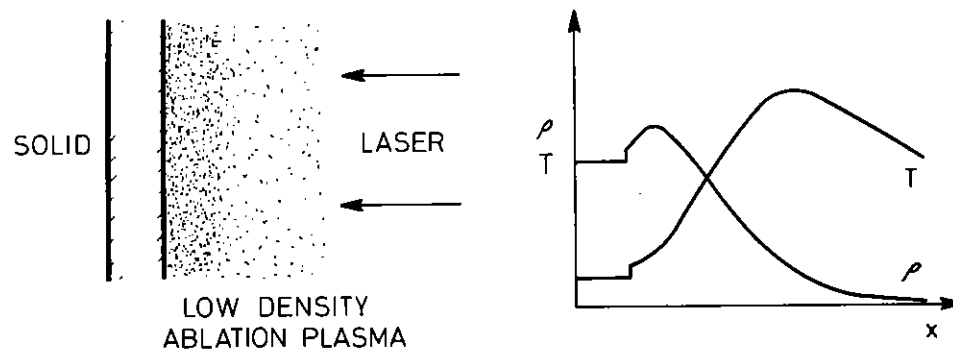
It is frequently valuable to have a simpler physical model which is analytically tractable and which reveals the basic behaviour of the target, even though the scope and accuracy of the simple model is less than that of the computer model. Various analytic models have been presented for laser driven ablation (8.13), (8.17), (8.18), (8.19) and in general they have become simpler and more transparent as time has passed.

8.2.4(a) Laser Driven Ablation

Laser driven ablation produces a low density plasma around the target within a very short space of time. The laser radiation is unable to penetrate beyond its critical density ρ_c ($\rho_c = 3.3 \times 10^{-3} \lambda^{-2}$; λ = laser wavelength in μm) and so the heating is confined to the plasma atmosphere (see Figure 8.03).

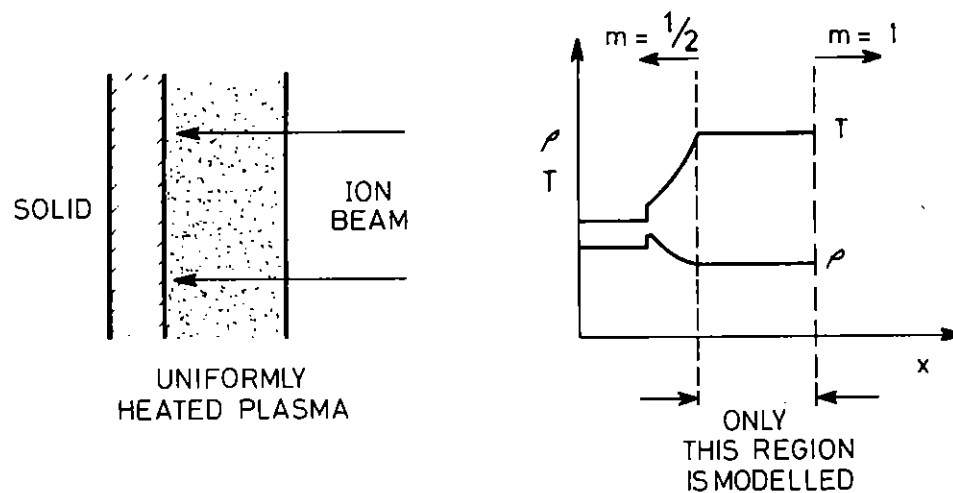
Analysis of the steady state continuity and momentum equations shows that the system must have a sonic point close to the critical density. The effective ablation pressure p_A is given in terms of the absorbed irradiance I and the sonic density ρ_1 as

$$p_A = 2 (3)^{-2/3} I^{2/3} \rho_1^{1/3}$$



LASER IRRADIATED SOLID

Fig. 8.03 Schematic representation of density and temperature profiles for a laser irradiated solid.



ION BEAM IRRADIATED SOLID

Fig. 8.04 Schematic representation of density and temperature in the idealised model of an ion beam heated target.

Calculations by Manheimer, Colombant and Gardner (8.13), including a more careful treatment of the rarefaction wave and energy balance give

$$P_A = 11 I_{14}^{2/3} \lambda^{-2/3} \text{ Mbar}$$

where I_{14} is the absorbed irradiance in units of $10^{14} \text{ W cm}^{-2}$, and λ is the laser wavelength in μm .

8.2.4(b) Ion Beam Heating

In contrast to the case of laser driven ablation the targets heated by ion beams do not approach a steady state within the timescales of interest. The particle beams have a given "range" $R \text{ gm cm}^{-2}$ which is only weakly dependent on the density and temperature of the target (and this dependence will be ignored). The energy deposition rate is generally a decreasing function of the projectile particle energy E resulting in the "Bragg peak" where the energy deposition is a maximum near the end of the range. Finite temperature effects in the target material (8.20), (8.14), and change of effective charge with velocity for heavy ions (8.15) both act to remove the Bragg peak and we will assume for simplicity that the specific energy deposition rate ie $\frac{\partial E}{\partial(\rho x)}$ is constant.

In response to this heating by the particle beam the target will expand into the vacuum on one side, and drive a shock wave into the solid on the other side (see Figure 8.04). A simple model of the hydrodynamic behaviour would therefore be to have one surface expanding, at a Mach number $M = 1$, and the other at $M = M'$ corresponding to the piston velocity behind a strong shock in a perfect gas. The material density in the target will not in general be constant but we will make this assumption in the spirit of a simple model, and the validity of the approximation will be borne out later by the agreement with more sophisticated calculations.

If we consider a slab of material of density ρ and thickness d such that $\rho d = R$; the range of the beam particles, then

$$d(t) = d_0 + (1 + M') \int_0^t C_s(t') dt'$$

The Mach number of the "piston" is initially $M' = 0.671$ but this reduces as the heated region expands and we take an average value $M' = 0.5$.

The energy equation relates the rate of rise in temperature to the energy input and the work done in hydrodynamic expansion

$$\frac{3}{2} \frac{k}{M_1} \frac{dT}{dt} = \frac{I}{R} - \frac{1.5 p C_s}{R} \quad (8)$$

where k is Boltzmann's constant, M_1 the average mass per particle of the target material and the other symbols have been defined previously.

$$\text{If we put } \theta = C_s^2 = \frac{p}{\rho} \frac{kT}{M_1}$$

$$\text{then } \frac{3}{2} \frac{d\theta}{dt} = \frac{I}{R} - \frac{1.5 \rho \theta^{3/2}}{R}$$

$$\text{and } \rho = \frac{R}{d + 1.5 \int_0^t \theta^{1/2} dt'}$$

By scaling the variables we may reduce this to a single dimensionless equation:

$$\theta = \frac{\theta R^{3/2}}{(I d_0)^{2/3}}$$

$$\tau = \frac{t I^{1/3}}{d_0^{2/3} R^{1/3}}$$

$$\text{giving } \frac{3}{2} \frac{d\theta}{d\tau} = 1 - \frac{1.5 \theta^{3/2}}{1 + 1.5 \int_0^\tau \theta^{1/2} d\tau'} \quad (9)$$

This integro-differential equation may be integrated numerically to give

$\theta(\tau)$ subject to $\theta(0) = 0$. From this we may obtain the scaled density and pressure as functions of time

$$\chi = \frac{\rho}{\rho_0} = (1 + 1.5 \int_0^\tau \theta^{1/2} d\tau')^{-1}$$

$$\text{and } P = \theta \times \chi$$

These numerical solutions are plotted in Figure 8.05. The temperature increases steadily in time while the density decays and the pressure passes through a very broad maximum $P = 0.32$ for $1 \lesssim \tau \lesssim 2$. It is reasonable to take this quasi-steady pressure as the "ablation pressure" generated by the ion beam. Transforming to the physical variable p , T and t gives, at $\tau = 1$:

$$\begin{aligned} p &= 32 I_{14}^{2/3} \rho_0^{1/3} \text{ Mbar} \\ T &= 114 I_{14}^{2/3} \rho_0^{-1/3} \text{ eV} \\ t &= 10^{-7} R I_{14}^{-1/3} \rho_0^{-2/3} \text{ sec} \end{aligned} \quad (10)$$

For the ion beam heated targets the ablation pressure shows the same $I^{2/3}$ dependence as for laser targets but ρ_1 ; the sonic density of the laser targets, has been replaced by ρ_0 ; the solid density. The ion beam generates much more pressure for a given intensity ie it has a higher hydrodynamic efficiency. It is of considerable interest to note that the peak pressure generated by the ion beam is independent of the range R of the charged particles. However the time scale over which the pressure persists is dependent on the range. For a given pulse duration, if R is too small the pressure will be transient and will decay during the pulse, if R is too large the pressure will not reach its maximum during the drive pulse.

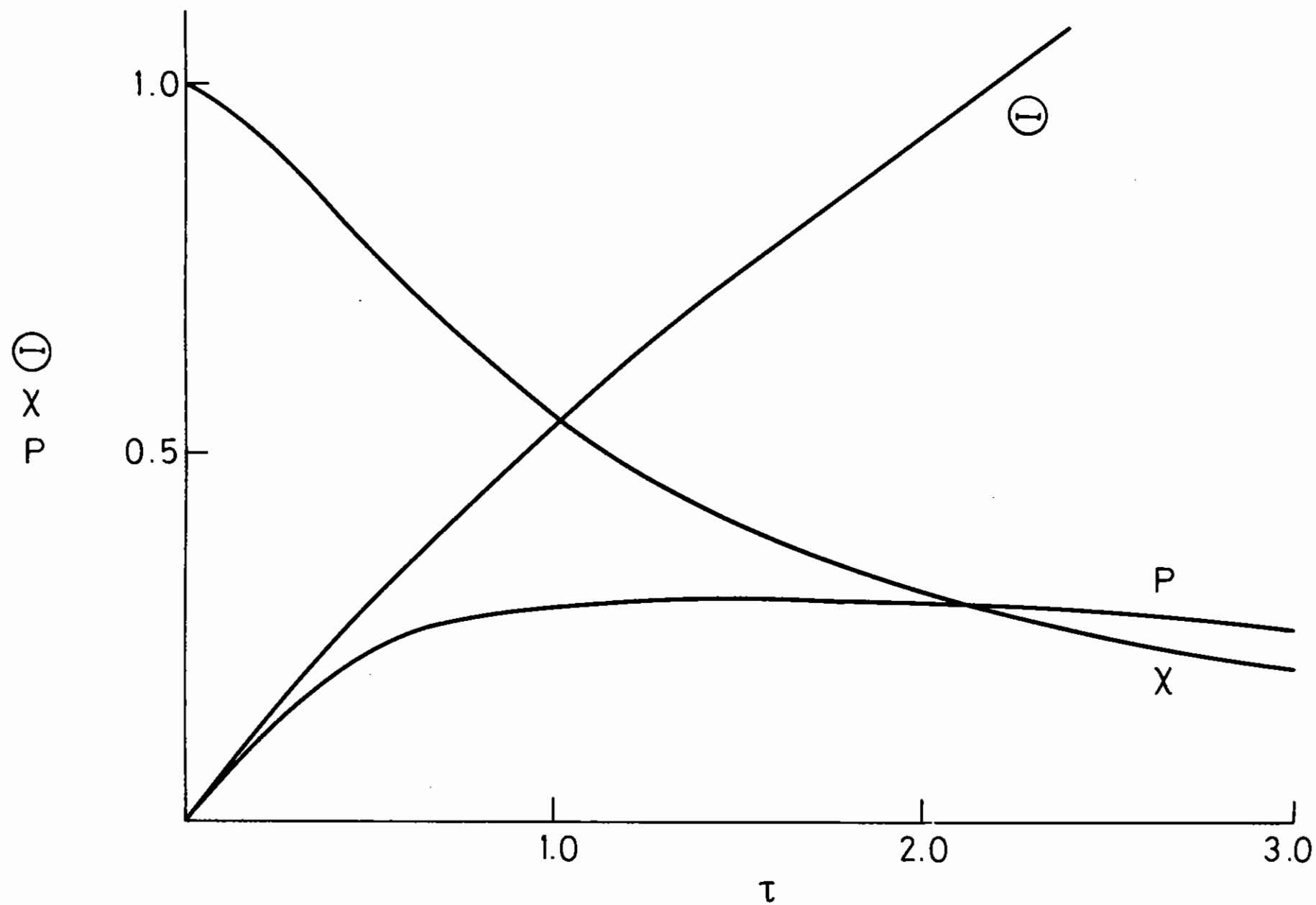


Fig. 8.05 Solutions for the scaled temperature Θ , density λ , and pressure p , as a function of the scaled time τ .

8.2.4(c) Comparison with computer simulations

The MEDUSA computer code has been used to simulate ion beam heated targets (8.16) with the energy deposition function either having a Bragg peak or being constant as assumed here. The two sets of numerical results showed only a small difference but for consistency the comparison will be made with the uniform deposition simulations. The simulations were performed for I varying from $10^{12} \text{ W cm}^{-2}$ to $10^{15} \text{ W cm}^{-2}$ in decade steps, and R taking the values 0.01, 0.03, 0.1, 0.3 gm cm^{-2} . The pressures quoted are all measured at $t = 5 \text{ ns}$ and so do not always correspond to the maximum pressure in the simulation runs. The MEDUSA code included thermal conduction and real equation of state effects as well as a solution of the hydrodynamic equations but did not include radiative energy losses. In Figure 8.06 the pressures measured in the simulations are compared with the results of equation (8) and the agreement is obviously very good. Some of the simulation data lies above the analytic prediction of maximum pressure and this corresponds to the degree to which the target density and pressure do not remain uniform, as assumed in the simple model.

8.2.4(d) Limitations of the model

The results of the previous section show that the simple analytic model is in good agreement with a more sophisticated computer model over a wide range of values of I and R . However the particular computer model used did not include radiative energy loss and this is the major source of error. Ion beams tend to heat a large depth of the target material with the result that optically thick black body radiation is emitted. The effective black body temperature will depend on the details of the radiation transport and the temperature profile near the surface. A useful check on the simple model is to require that the black body radiation flux at the (uniform) temperature T is less than the particle beam energy flux I i.e. $\sigma T^4 < I$. If we take the temperature at $\tau = 1$ then

$$T = 114 I^{2/3} \rho_0^{-2/3} \text{ (eV)} \quad (11)$$

and the condition $\sigma T^4 < I$ is equivalent to

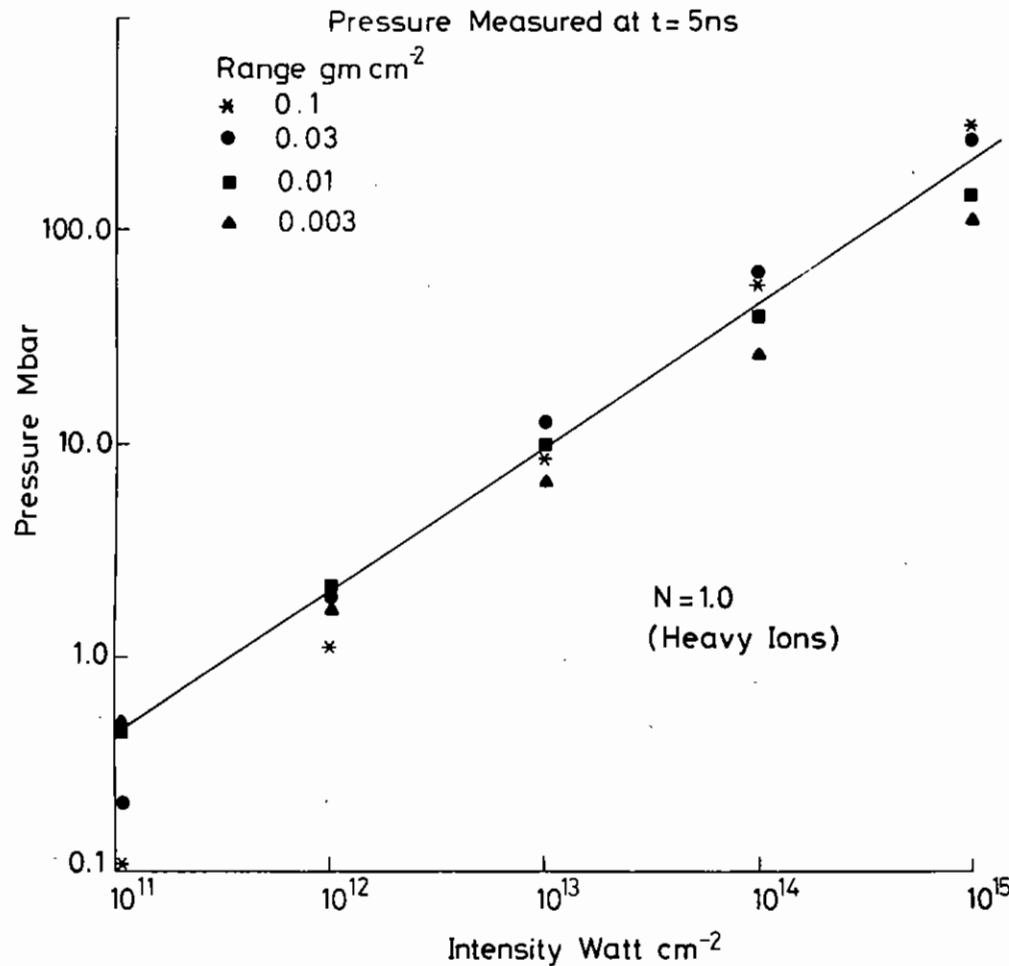


Fig. 8.06 Comparison of the analytic model (full line) with Medusa computer calculations, for the ablation pressure as a function of beam intensity.

$$I_{14} < 0.3 \rho_0^{8/5} \quad (12)$$

or $I < 1.4 \times 10^{14} \text{ W cm}^{-2}$ for $\rho_0 = 2.7 \text{ gm cm}^{-3}$ (aluminium). Since the black body losses vary so rapidly with temperature even a small reduction in I below this limit makes the radiated losses insignificant. The maximum pressure that can be generated without excessive radiative losses is given by equations (10) and (12) as $p_{\text{max}} = 10 \rho^{7/5} \text{ Mbar}$.

8.2.4(e) Conclusions

The quasi steady pressure generated by an ion beam does not depend on the range of the beam particles but the time scale of the process does. For ICF pellet applications with $I \approx 10^{14} \text{ W cm}^{-2}$ and $\tau \approx 10^{-8} \text{ sec}$ a particle range of $0.1 - 0.2 \text{ g cm}^{-2}$ is appropriate. A pressure of more than 100 Mbar may be generated without the radiated energy losses exceeding 30% of the input energy. Targets with layers of different density and chemical composition may of course perform better than these simple considerations suggest.

R G Evans (RAL)

8.2.5 Soft X-ray Preheat as a Constraint on Ablative Implosions

8.2.5(a) Introduction

In an ablative implosion the ablation pressure P_a generated at the outer surface of a spherical shell target creates a much higher pressure P_f at the centre of the shell at the time of implosion. This pressure P_f is proportional to the maximum momentum flux density $\rho_s v_s^2$ of the imploding shell (8.21). For a simple unshaped laser pulse the density of the imploding shell ρ_s at peak velocity is that due to compression by a shock of pressure P_a . Preheating (by hot electrons or soft X-rays) raises the internal energy of the imploding shell and its hydrodynamic response is to decompress in order to maintain a maximum pressure equal to the applied ablation pressure P_a . This preheat induced decompression

reduces the pressure P_f and the density ρ_f at the time of implosion (a) because ρ_s is reduced (b) because v_s is reduced through earlier closure of the void inside the shell and therefore earlier onset of deceleration (8.22); (c) because the spherical convergence enhancement of P_f (8.21) is reduced in a lower density implosion.

An important constraint for the optimisation of ablative implosions is imposed by the need to make decompression by preheat negligible. In general this will be possible if the size and therefore wall thickness of the target are sufficiently large. Since the laser energy required to drive the target scales as (radius)³, there are major advantages in using small targets to study implosion physics but the minimum laser energy and target size are determined by the tolerable level of preheating.

These scaling issues have been discussed in some detail for preheating by hot electrons (8.21) but with the advent of ablative implosions driven by short wavelength lasers ($\lambda < 0.5 \mu\text{m}$) there are circumstances where preheating by soft X-rays may be more important.

Numerical simulation of the hydrodynamics of implosions with a full treatment of the relevant atomic physics and radiation transport gives the most accurate description of the effects of soft X-ray preheat (see section 8.4.9 *ibid*). It is however complex and gives little physical insight or understanding of scaling behaviour. A simpler analytic approach may therefore be useful.

8.2.5(b) Model description of Soft X-ray Preheating

The source of soft X-ray preheating is the hot plasma at the ablation front. The structure of the plasma (Figure 8.07) is such that radiation travelling outwards from the target encounters a negligible mass per unit area of target material relative to that encountered in travelling to the inner half of the shell, (where we require an estimate of the preheating). Preheating in the inner half of the shell is due to photons with a range of the order of the shell thickness. Photons of this energy are therefore not significantly absorbed in traversing the relatively small amount of

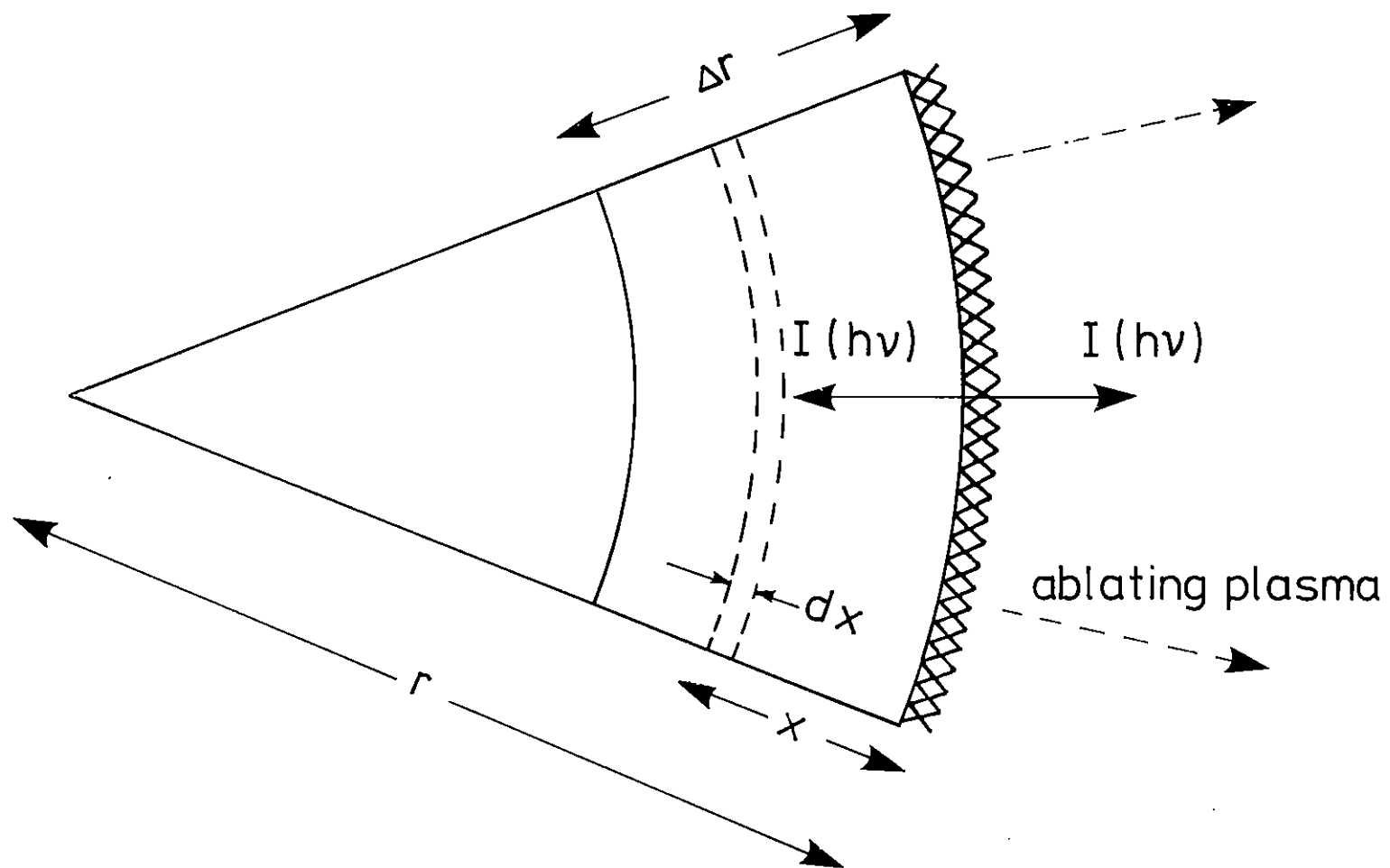


Fig. 8.07 Geometry and notation used in the analysis of preheat.

target material which they encounter in escaping radially outwards from the target. It follows that the outward radiation flux density $I(h\nu)$ from the target can be used to assess preheating by assuming an isotropic source with a similar flux density $I(h\nu)$ directed radially inwards.

$I(h\nu)$ can be measured experimentally (section 7.5 *ibid*). The radiation is mainly free free and free bound continuum whose spectrum can be locally approximated by $I(h\nu) \propto \exp(-h\nu/kT)$, where kT is the source temperature. kT increases with $h\nu$ since higher energy photons originate from hotter less dense plasma further from the target surface (see eg section 7.5). The exponential form is characteristic of an optically thin source and for photon energies sufficient to penetrate to the inner half of the shell the ablation plasma is optically thin.

The specific preheating rate $\dot{\epsilon}$ at depth x in the shell can be written

$$\dot{\epsilon} = \int_0^{\infty} I(h\nu) \exp\left[-\int_0^x a(h\nu, x') \rho(x') dx'\right] a(h\nu, x) d(h\nu) \quad (1)$$

where $a(h\nu, x)$ is the X-ray mass absorption coefficient and $\rho(x)$ the plasma density.

This can be simplified to

$$\dot{\epsilon} \approx \int_0^{\infty} I(h\nu) \exp[-a(h\nu) \bar{\rho} x] a(h\nu) d(h\nu) \quad (2)$$

where $\bar{\rho} x$ is the depth in units of mass per unit area and $a(h\nu)$ is constant for a given frequency. [The latter is a reasonable approximation when absorption is mainly due to K shell electrons (eg C in polymer and O in glass) since the K shell typically remains unionized for temperature produced by shock and radiation preheat.]

The temperature rise kT_p due to preheating can be readily calculated for an ideal fully ionized plasma ($A \approx 2Z \gg 1$) viz

$$kT_p = 4/3 m_p \dot{\epsilon} \tau \quad (3)$$

where m_p is the proton mass and τ the duration of the preheating. [In practice real equation of state effects are important and (3) underestimates the temperature by as much as a factor of 2.]

This temperature rise kT_p can be compared with that produced by shock heating kT_s , when an ablation pressure P_a is applied suddenly to a target. In the same ideal plasma limit we have

$$kT_s / 1 \text{ eV} = 0.5 \frac{P_a / 1 \text{ Mbar}}{\rho / 1 \text{ gcm}^{-3}} \quad (4)$$

where ρ is the density of the shell.

It is convenient, when evaluating kT_p from experimental data on $I(h\nu)$, to know which spectral regions are important. This can be achieved by introducing the 'cut off' photon energy $h\nu_c$ defined by

$$a(h\nu_c) \bar{\rho} x = 1 \quad (5)$$

and by expressing the spectrum $I(h\nu)$ in the vicinity of $h\nu_c$ as

$$I(h\nu) = I(h\nu_c) \exp(h\nu_c/kT) \exp(-h\nu/kT) \quad (6)$$

and the mass absorption coefficient $a(h\nu)$ as,

$$a(h\nu) = a(h\nu_c) (h\nu_c/h\nu)^3 \quad (7)$$

Then with the variable $y = h\nu/h\nu_c$ and introducing $b = kT/h\nu_c$ we have from equations (2), (3), (5), (6) and (7)

$$kT_p / 1 \text{ eV} = 2 \times 10^{-30} (\tau / 1 \text{ ns}) (I(h\nu_c) / (1 \text{ keV keV}^{-1} \text{ cm}^{-2} \text{ sec}^{-1})) x \dots$$

$$(h\nu_c / 1 \text{ keV}) (\bar{\rho} x / 1 \text{ gcm}^{-2})^{-1} x \dots \quad (8)$$

$$\int_0^{\infty} \exp(1/b) \exp(-y/b) \exp(-y^{-3}) y^{-3} dy$$

The frequency dependence of the integral in (8) is shown in Figure 8.08 for various values of the parameter b . In all cases it peaks at $h\nu = h\nu_c$ and has a fairly narrow range about $h\nu_c$ showing that preheating is due to photons with $h\nu = h\nu_c$.

The value of the integral $F(b)$ is a function of the parameter b as shown in Figure 8.09, and it is clear that it is essentially independent of b , implying that preheating is mainly determined by $I(h\nu_c)$ and that the slope of the spectrum is relatively unimportant. Measurement of $I(h\nu)$ in the vicinity of $h\nu_c$ is therefore all that is required to estimate the preheating.

8.2.5(c) Preheat in ablative implosions

In an ablative implosion the implosion time

$$\tau_I = r \left(\frac{\rho \Delta r}{P_a} \right)^{\frac{1}{2}} \quad (9)$$

where Δr is the thickness of the shell so that at the mid point of the shell from equations (8) and (9)

$$kT_p / 1 \text{ eV} = 1.2 \times 10^{-27} \left(\frac{r}{\Delta r} \right)^{\frac{1}{2}} \left(\frac{1}{\rho / 1 \text{ g cm}^{-3} P_a / 1 \text{ Mbar}} \right)^{\frac{1}{2}} \times \dots \quad (10)$$

$$\dots (I(h\nu_c) / 1 \text{ keV keV}^{-1} \text{ cm}^{-2} \text{ sec}^{-1}) (h\nu_c / 1 \text{ keV})$$

It is also useful to consider the relative importance of shock and radiative preheating which can be written from equations (4) and (10) as

$$kT_p / kT_s = 2.4 \times 10^{-27} \left(\frac{r}{\Delta r} \right)^{\frac{1}{2}} \frac{(\rho)^{\frac{1}{2}}}{(P_a)^{3/2}} I(h\nu_c) (h\nu_c) \quad (11)$$

using the same units as equation (10).

A criterion for a satisfactory ablative implosion is $kT_p / kT_s < 1$.

Since the aspect ratio $r/\Delta r$ is fixed at about 10 by the need to control Rayleigh Taylor instability, the free parameters are ablation pressure P_a , density ρ and atomic composition, cut off energy $h\nu_c$ and $I(h\nu_c)$.

P_a is a function of absorbed irradiance I_a and for irradiation at $\lambda = 0.53 \mu\text{m}$ we can put (8.23).

$$P_a / 1 \text{ Mbar} \approx 2(I_a / 10^{13} \text{ W cm}^{-2})^{0.86} \quad (12)$$

$I(h\nu_c)$ is a function of irradiance I_a and atomic composition of the ablation plasma. If we assume as a rough approximation that for a given material the conversion efficiency to X-ray emission is constant, then $I(h\nu_c) \propto I_a$ which with equation (11) and (12) implies that $kT_p / kT_s \propto I_a^{-0.3}$. The weak scaling suggests that irradiance is a relatively unimportant parameter in determining preheat in an implosion.

The choice of atomic composition of the shell and the related density ρ are influenced by the required implosion velocity. If high velocity and therefore high temperature in the compression are needed (to produce thermonuclear and X-ray emission for example) then for a given $r/\Delta r$ a low value of ρ is advantageous, $v_s^2 \propto (P_a / \rho)(r/\Delta r)$. Conversely higher ρ gives a lower temperature higher density compression (8.21).

In the case of a homogeneous shell the remaining free parameter is $h\nu_c$ which is determined by Δr and therefore by the overall target size r (for fixed $r/\Delta r$). The scaling is $h\nu_c \propto \Delta r^{1/3}$ (equations (5) and (7)). This weak scaling is offset by the exponential form of $I(h\nu_c) \propto \exp(-h\nu_c / kT)$, so that increasing the target size has a strong effect on preheat as shown in more detail below.

Inhomogeneous shells may be advantageous. An obvious case is the reduction of $I(h\nu_c)$ obtained with a low Z ablation layer on a shell of higher Z . Another possibility is the inclusion of a layer of high mass

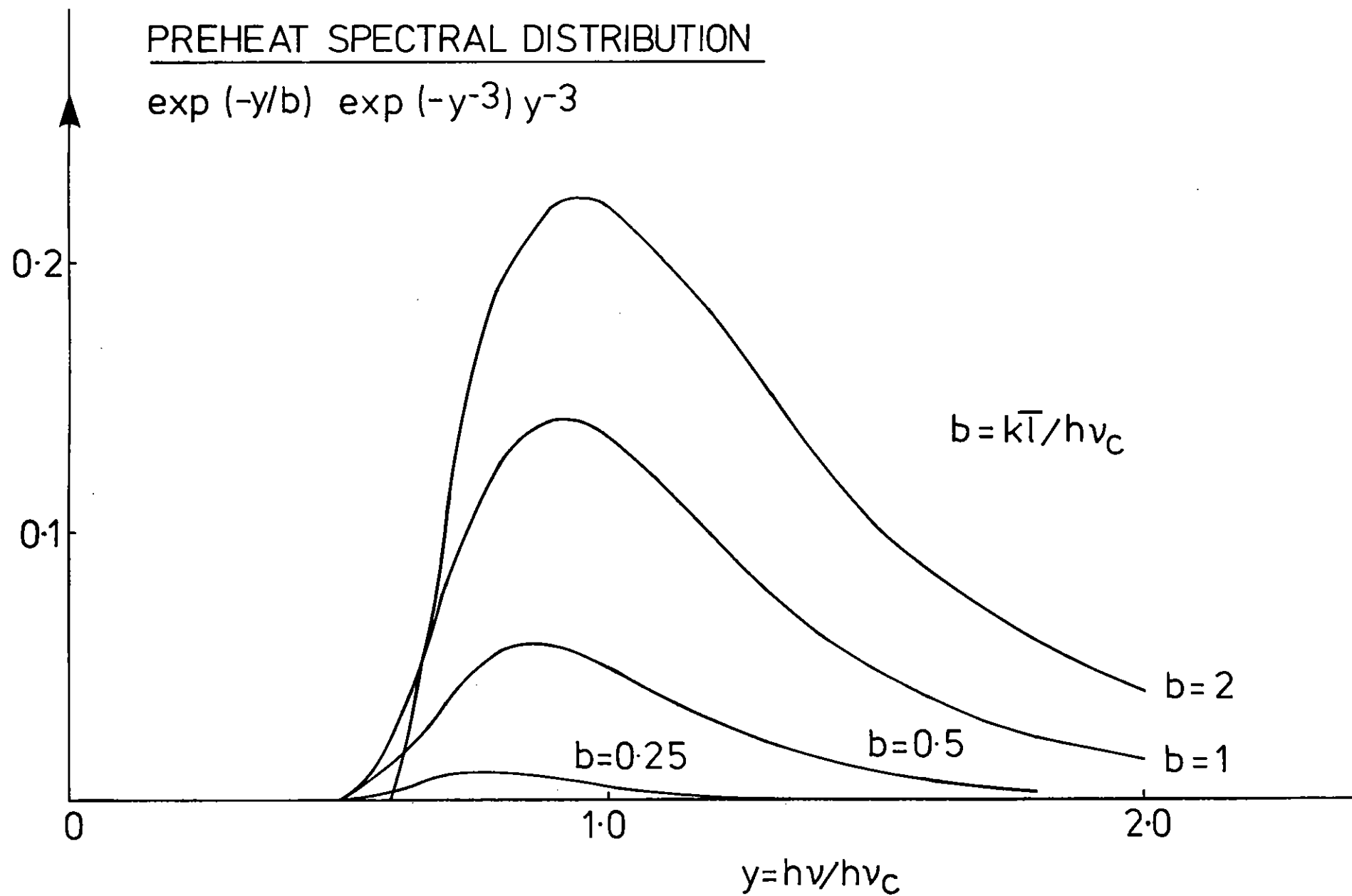


Fig. 8.08 Spectral distribution of preheat.

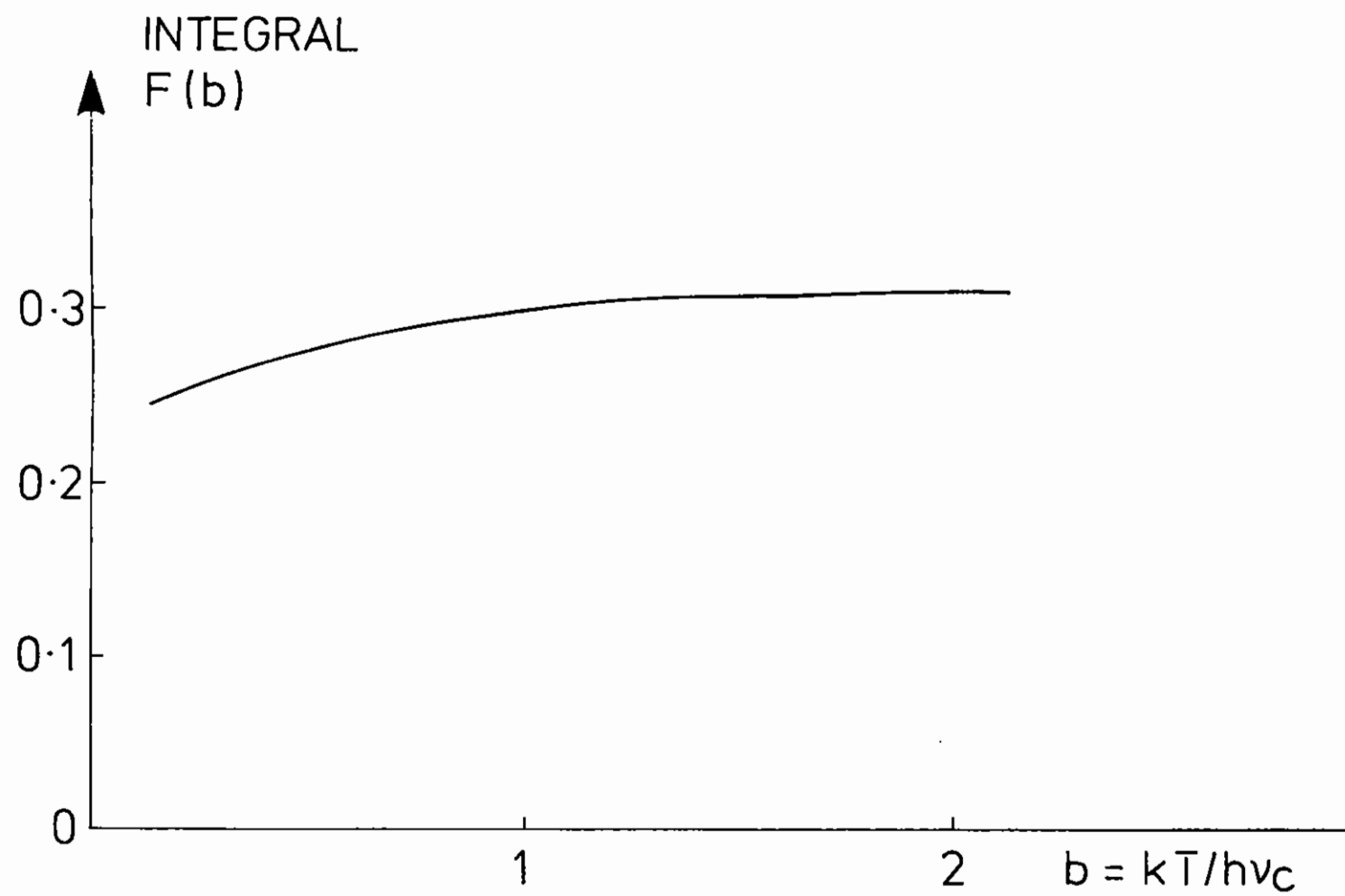


Fig. 8.09 Graph of the integral $F(b)$.

absorption coefficient as a preheat barrier. The latter is less useful than it might at first sight appear because of the relatively small difference ($\sim 3 \times$) between the highest available mass absorption coefficients and that of glass. The disadvantages of dense RT unstable layers may outweigh the advantage of the small increase in preheat screening.

8.2.5(d) Estimating Preheat

Experimental or theoretical data giving $I(h\nu_c)$ are required for quantitative estimates of preheat using the above model. Section 7.5 *ibid* compares some experimental and calculated results and other data are available in the literature.

An interesting example is a simple polymer shell of low aspect ratio. X-radiography studies of the implosion of targets of this kind is discussed in section 6.2 *ibid* for shells with $r = 60 \mu\text{m}$, $\Delta r = 10 \mu\text{m}$, $I_a = 3 \times 10^{13} \text{ W cm}^{-2}$, $\lambda = 0.53 \mu\text{m}$, $\tau = 1.3 \text{ ns}$.

The mass absorption coefficient for a typical polymer $(\text{CH}_2)_n$ is shown in Figure 8.10, together with a compilation of experimental and theoretical data giving $I(h\nu)$ for an absorbed irradiance of $4 \times 10^{13} \text{ W cm}^{-2}$, the latter taken from section 7.5 *ibid*. For the above example of a polymer shell we have $h\nu_c = 1050 \text{ eV}$ for penetration to a depth of $5 \mu\text{m}$ (half way through the shell). The corresponding value of $I(h\nu_c)$, scaled to the experimental absorbed irradiance is (from Figure 8.10) $1.5 \times 10^{27} \text{ keV keV}^{-1} \text{ cm}^{-2} \text{ sec}^{-1}$. Hence from equation (8) with $F(b) = 0.3$ we have $kT_p \sim 2.3 \text{ eV}$. The situation is complicated a little by preheat in the transmission window below the C K edge at 285 eV (Figure 8.10). Extrapolating the mass absorption coefficient below the edge gives $h\nu_c = 298 \text{ eV}$. It follows that about half the frequency range of the preheating band in Figure 8.08 ($0 < h\nu/h\nu_c < 0.95$) will contribute to preheating. A rough estimate of $I(h\nu_c)$ is given by the calculated data in Figure 8.10 viz $6 \times 10^{27} \text{ keV keV}^{-1} \text{ cm}^{-2} \text{ sec}^{-1}$, and hence an estimate of the preheating in the transmission window which is $kT_p \sim 0.6 \text{ eV}$. The total preheat is thus estimated at 3 eV or $2 \times 10^5 \text{ J/g}$ compared to the

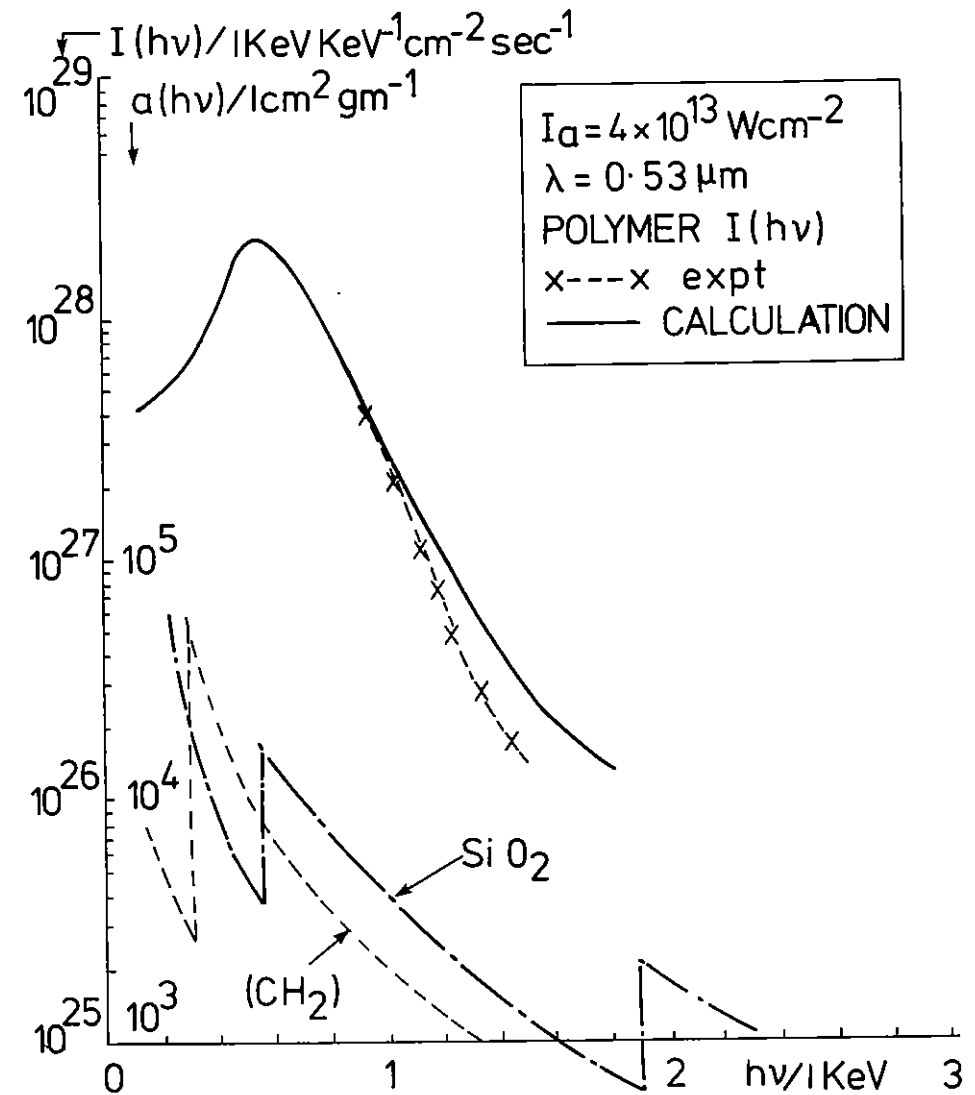


Fig. 8.10 Mass absorption coefficient.

above C				below C		$(kT_{p1} + kT_{p2})/kT_s$
		K edge		K edge		
r	Δr	$h\nu_c$	kT_{p1}	$h\nu_c$	kT_{p2}	
μm	μm	eV	eV	eV	eV	
60	10	1050	2.3	288	0.6	1.4
120	20	1320	0.3	380	0.3	0.3

Table 8.01 Polymer shell $I_a = 3 \times 10^{13} \text{ W cm}^{-2}$ $\lambda = 0.53 \mu\text{m}$

r	Δr	τ	Si window		O window		$(kT_{p1} + kT_{p2})/kT_s$
			$h\nu_c$ eV	kT_{p1} eV	$h\nu_c$ eV	kT_{p2} eV	
μm	μm	ns	eV	eV	eV	eV	
80	1.0	0.3	770	120	417	30	44.0
80	5.0	0.7	1320	3.4	713	18	6.0

Table 8.02 Glass shell $I_a = 1.2 \times 10^{14} \text{ W cm}^{-2}$ $\lambda = 0.53 \mu\text{m}$

shock preheat value equation (4) of 2 eV or $1.3 \times 10^5 \text{ J/g}$. The target in this example is therefore close to the limit for satisfactory ablative compression since $kT_p/kT_s \sim 1.5$ (see Table 8.01).

A check on the accuracy of the simple model was made by comparing the above estimate with detailed numerical modelling using the method described in section 8.4.7 *ibid*. This indicated a total preheating (shock + radiation) of 5 eV for the shell at peak velocity $v_f = 4 \times 10^6 \text{ cm sec}^{-1}$ with density ρ_s of 2.5 g cm^{-3} . The agreement with the model calculated preheat (shock + radiation) of 4.9 eV is very good.

The beneficial effect of an increase in scale size and therefore wall thickness is shown by repeating the calculation for a 2 x larger target with a 2 x longer implosion time as in Table 8.01. Preheating is seen to be reduced 2 fold.

A similar calculation for glass shell targets involves considering preheating in the spectral region below the Si K edge where O K shell absorption is dominant, and also an additional contribution from below the O K edge where O and Si L shell absorption predominate (see Figure 8.10). A glass shell has about a 4 x greater soft X-ray source intensity (neglecting line radiation) than a polymer shell (section 7.5 *ibid*). With this assumption Figure 8.10 can be used to estimate $I(h\nu)$.

An interesting example to consider is the hot ablative implosion of high aspect ratio glass shells described in section 6.4 *ibid* where $r = 80 \mu\text{m}$, $\Delta r = 1.0 \mu\text{m}$, $I_a = 1.2 \times 10^{14}$ and $\tau_L = 0.3 \text{ ns}$. Table 8.02 shows calculations of preheat by the above method for this target. The radiation preheat level is clearly large relative to shock preheating ($kT_p/kT_s = 44$) indicating an unsatisfactory choice of target for an ablative implosion.

A check against a full numerical simulation as described in section 8.4.7 *ibid* was made for this example also, and indicated $kT_p + kT_s = 180 \text{ eV}$ at peak velocity $v_s = 8 \times 10^6 \text{ cm sec}^{-1}$ with $\rho_s = 0.3 \text{ g cm}^{-3}$. There is again fair agreement with the model result $kT_p + kT_s = 150 \text{ eV}$ for a

situation dominated by preheating. A thicker wall $\Delta r = 6 \mu\text{m}$ with the same radius with correspondingly longer (0.7 ns) implosion time is also analysed in Table 8.02 to illustrate the 7 x improvement in preheat screening. [The 0 window contribution in this estimate is uncertain within a factor of a few time since experimental data on $I(h\nu)$ are not available and $a(h\nu)$ is sensitive to preheat.]

8.2.5(e) Conclusion

A simple model description gives a rapid and useful indication of the minimum target size needed for ablative implosion of simple polymer and glass shells in order that soft X-ray preheat does not degrade the compression. The preheat is seen to be determined by the soft X-ray intensity at the cut off frequency $I(h\nu_c)$ indicating, for minimum sized targets, a need for experimental data in two spectral regions, viz below the 0 and C K edges and in the 1 to 2 keV region. A useful conclusion is that the relative importance of radiative and shock preheating is insensitive to the irradiance.

M H Key (RAL)

8.3 Laser Plasma Interactions

8.3.1 Oscillations due to feedback in interaction processes

Emissions from laser produced plasmas may show oscillatory behaviour, a notable example being that at $3/2 \omega_0$ (8.24). Here we consider, in a general way, the conditions under which oscillations may be expected to occur as a result of feedback effects with time delay, then look at some particular processes.

We suppose that the observed emission, E say, is a function of some plasma parameter X,

$$E = F(X) \quad (1)$$

Now suppose that there is some mechanism by which the process giving rise to E affects X. To take the example of $3/2 \omega_0$ emission; if this is supposed to be due to two-plasmon decay near the quarter-critical surface, then the growth rate, and presumably the saturated intensity, depends on the gradient length scale. This latter quantity is, however, affected by the ponderomotive force generated by the decay waves, and so will depend on the wave intensity generated by the instability. Our object is to consider such feedback loops and, supposing that there is some time delay in a non-stationary system, to investigate the stability of stationary states.

In a stationary state we take

$$X = G(E) \quad (2)$$

so that a self consistent equilibrium will satisfy (1) and (2) simultaneously. Now, in a time-dependent system we suppose that

$$X(t) = \int_0^\infty g(\tau, E(t-\tau)) d\tau \quad (3)$$

representing a fairly general form of dependence of X on the past history of E. If E_0 and X_0 are the equilibrium values,

$$X_0 = \int_0^\infty g(\tau, E_0) d\tau = G(E_0) \quad (4)$$

Now we examine the stability of the equilibrium by letting

$$E = E_0 + E_1(t)$$

$$X = X_0 + X_1(t)$$

and linearising (3), to get

$$X_0 + X_1(t) = G(E_0) + \int_0^\infty \frac{\partial g(\tau, E_0)}{\partial E_0} E_1(t-\tau) d\tau$$

Using equation (1) to express E_1 in terms of X_1 and supposing X_1 to go as e^{pt} we get

$$1 = \frac{\partial F(X_0)}{\partial X_0} \frac{\partial g(\tau, E_0)}{\partial E_0} e^{-p\tau} d\tau \quad (4)$$

Stability of the system depends on whether (4) has solutions for p with a positive real part.

If $\frac{\partial F}{\partial X_0}$ and $\frac{\partial G}{\partial E_0}$ have the same sign it is easy to see that if $\frac{\partial F}{\partial X_0} \frac{\partial G}{\partial E_0} > 1$ the right hand side of (4) is greater than 1 for $p = 0$ and tends to zero as $p \rightarrow \infty$, so there must be a solution of (4) with positive p and the equilibrium is unstable. Unfortunately the derivatives are generally of opposite sign, in which case there is no such simple criterion and the time-dependence of g is important. If, however, we make the approximation that $g(\tau, E_0)$ is constant in the interval $\tau_1 < \tau < \tau_2$, and zero otherwise, then we get

$$1 = \frac{\partial F}{\partial X_0} \frac{\partial G}{\partial E_0} \frac{e^{-p\tau_1} - e^{-p\tau_2}}{p(\tau_2 - \tau_1)}$$

or, if $\frac{\partial F}{\partial X_0} \frac{\partial G}{\partial E_0} = -a$, $p' = p\tau_1$ and $T = \tau_2/\tau_1$,

$$- \frac{1}{a} = \frac{e^{-p'} - e^{-p'T}}{p'(T-1)} \quad (5)$$

The value of a at which there is a transition from stability to instability of the equilibrium is that which makes p' purely imaginary, and it can be shown from (5) that the smallest such value is

$$a = \frac{\pi(T-1)}{T+1} \operatorname{cosec} \left(\frac{2\pi}{T+1} \right) \quad (6)$$

This threshold value of a , above which the equilibrium is unstable tends to 1 as T tends to 1 (corresponding to a δ -function for the time-dependence of g) and increases with T .

Our conclusion then is that information about the equilibrium of the system can be obtained from a steady-state analysis in which it is only

necessary to evaluate the derivatives $\frac{\partial F}{\partial X_0}$ and $\frac{\partial G}{\partial E_0}$. Instability will occur if $\frac{\partial F}{\partial X_0} \frac{\partial G}{\partial E_0}$ is greater than some critical value. The effect on this value of the time-dependence of the feedback effect, given by (6) for a particular form of time-dependence, is not very strong. For example if a in (6) is greater than about 5 instability will occur for values of T up to about 10, allowing considerable latitude in the time-dependence of g . We may also note that the critical value of a does not depend on the absolute time scale over which the feedback occurs, but only on the shape of the time-dependence.

To illustrate this theory we consider two-plasmon decay. It has been suggested that the observed oscillations in $3/2 \omega_0$ emission are due to the fact that the growth in oscillating field amplitude causes a steepening of the density gradient near the critical surface which switches off the instability. The plasma then relaxes to a less steep gradient and the process recommences. The question which we consider is why this should happen, rather than simply the setting up of an equilibrium in which steepening of the density gradient is balanced by a self-consistent saturated amplitude of the instability, producing a steady profile like that which appears to occur at the critical surface. Using the above theory we investigate the stability of such an equilibrium. The time delay in the feedback process here may occur either because the plasma profile takes a finite time to adjust to a given ponderomotive force, or because the two-plasmon instability takes a finite time to reach a steady saturated amplitude when the profile is changed.

According to Liu (8.25), the maximum growth rate in an inhomogeneous plasma is

$$\gamma_{\max} = \frac{k_0 v_0}{2} - \frac{3v_e^2}{v_0 L}$$

with $v_0 = eE_0/m\omega_0$ and v_e the thermal velocity. We shall suppose that the saturated amplitude is simply proportional to the growth rate

$$E_{\text{sat}} = \beta \gamma_{\text{max}}$$

Then

$$\frac{dE_{\text{sat}}}{dL} = \frac{3v_e^2}{v_0 L^2} = \frac{3v_e^2}{v_0 L^2} \frac{E_{\text{sat}}}{\gamma_{\text{max}}} \quad (7)$$

We also need an equation giving the way in which the length scale depends on the saturated amplitude if the latter is given. To obtain this we shall simply assume that, as in the soliton solution of Chen and Liu (8.26), L is simply universally proportional to E_{sat} , ie $L = \delta/E_{\text{sat}}$ for some δ . Then

$$\frac{dL}{dE_{\text{sat}}} = - \frac{\delta}{E_{\text{sat}}^2} = - \frac{L}{E_{\text{sat}}} \quad (8)$$

The quantity on which stability depends is the product of (7) and (8), ie $-3v_e^2/v_0 L \gamma_{\text{max}}$. Now γ_{max} is less than $k_0 v_0/2$, so this is greater in magnitude than $\frac{e}{v_0^2} \frac{1}{k_0 L}$. For an irradiance of a few times 10^{14} W cm^2 and a temperature $\sim 60/k_0 L$. The equilibrium length scale depends on the saturated amplitude and would require a more detailed analysis. However, it seems likely that it would be such as to make the above quantity considerably greater than one. Thus, on the basis of a very simple calculation, we conclude that it is likely that the equilibrium profile set up near the quarter-critical density as a result of the ponderomotive force of the two-plasmon decay waves will be unstable and that oscillatory behaviour would indeed be expected.

Similar approximate calculations have been carried out to look at the possibility of oscillations of Brillouin scatter, assuming that its saturated level depends on electron temperature, which affects the damping of the ion sound wave, while the electron temperature depends on the amount of light reaching the critical surface, and hence on the level of Brillouin scatter. For most reasonable parameters and assumptions the indication is that this system is stable, which seems to be borne out by

observations of Brillouin scatter.

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8.3.2 The use of the finite element method in electrostatic particle codes

We describe a new application of the finite element method. We have used it in a two dimensional electrostatic particle-in-element (PIE) code, and have demonstrated its potential advantages by comparing collision and heating times and noise levels with those obtained from conventional codes using finite difference techniques. One consequence of using relatively small numbers of macroparticles in a particle simulation, is to raise the noise level. The most direct means of reducing noise levels is clearly to use more particles but this quickly makes the experiments prohibitively expensive. It is therefore worthwhile to look at all possible ways of combating the noise problem.

8.3.2(a) Method

In a finite element approach to the problem Poisson's equation is replaced by its functional analogue

$$\chi(\phi) = \int_R \left\{ \left(\frac{\partial \phi}{\partial x} \right)^2 + \left(\frac{\partial \phi}{\partial y} \right)^2 - 8\pi \rho \phi \right\} dx dy$$

The simulation region R is divided into rectangular elements and we write $\phi = N \phi_e^e$ for each element, where N are the shape functions and ϕ_e^e is the vector of nodal values. For each element R_e we therefore obtain the functional

$$\chi_e = \frac{1}{2} \phi_e^{eT} k_e^e \phi_e^e - \phi_e^{eT} F_e^e$$

where

$$k_e = 2 \int_{R_e} \left\{ \left(\frac{\partial N}{\partial x} \right)^T \left(\frac{\partial N}{\partial x} \right) + \left(\frac{\partial N}{\partial y} \right)^T \left(\frac{\partial N}{\partial y} \right) \right\} dx dy$$

and

$$F_e^e = 8\pi \int_{R_e} N^T \rho dx dy$$

If all functions are expressed in local coordinates (ξ, η) with origin at the lower left hand corner of an element of sides (a, b) , we have

$$k_e = 2 \int_0^1 \int_0^1 \left\{ \frac{b}{a} \left(\frac{\partial N}{\partial \xi} \right)^T \left(\frac{\partial N}{\partial \xi} \right) + \frac{a}{b} \left(\frac{\partial N}{\partial \eta} \right)^T \left(\frac{\partial N}{\partial \eta} \right) \right\} d\xi d\eta$$

and

$$F_e^e = 8\pi \int_0^1 \int_0^1 N^T(\xi, \eta) \rho ab d\xi d\eta$$

After assembling contributions from all elements and minimising with respect to each of the nodal parameters we obtain a set of linear equations

$$k\phi = F$$

which must be solved for the potential ϕ .

The electric field components are then

$$(E_x, E_y) = \left(-\frac{1}{a} \frac{\partial N}{\partial \xi} \phi^e, -\frac{1}{b} \frac{\partial N}{\partial \eta} \phi^e \right)$$

The equations governing the velocity $\tilde{v}_1$ and coordinates $\tilde{x}_1$ of the i th particle of mass m and charge e are

$$\frac{d\tilde{v}_1}{dt} = \frac{e}{m} \tilde{E}_1$$

$$\frac{d\tilde{x}_1}{dt} = \tilde{v}_1$$

where i will run over several thousand particles.

8.3.2(b) Collision and heating times

The lack of energy conservation is measured by the heating time of the simulation plasma which is defined to be the time for the average kinetic energy of an electron to increase by $\frac{1}{2}kT$.

The collision time for the simulation plasma is measured by first setting up a model in which the orbits of all particles would be straight lines in the Vlasov collisionless limit. The collisional effects can be measured by the deviations of the particle orbits from straight lines.

Suppose the initial velocity of a particle is $\tilde{v}_0$ and that after a time t its velocity is $\tilde{v}$. $\tilde{v}_\parallel(t) = |\tilde{v}| \cos \phi$, the component of velocity parallel to the original direction $h(t) = \frac{m}{2} (\tilde{v}^2 - \tilde{v}_0^2)$, the change in kinetic energy from its original value.

The following characteristic times may then be defined.

$$\text{Slowing down time, } \tau_S: \frac{d}{dt} \langle \tilde{v}_\parallel(t) \rangle = \frac{\langle \tilde{v}_\parallel(t) \rangle}{\tau_S} \text{ at } t = 0.$$

$$\text{Heating time, } \tau_H: \langle h(\tau_H) \rangle = \frac{1}{2}kT.$$

The brackets denote averaging over many particles.

Hockney also found that if N is the number density, H the mesh size and W the "effective particle width" then

$$\frac{\tau_S}{\tau_{pe}} = \frac{N(\lambda_D^2 + W^2)}{0.98} \quad (1)$$

where

$$\tau_{pe} = \frac{2\pi}{\omega_p}$$

and ω_p is the plasma frequency. He has concluded for the cloud-in-cell (CIC) method of charge sharing which is very often used in practice, that $W = H$.

Surprisingly in a more elementary scheme in which all the charge of a particle is associated with the nearest grid point (NGP) particles also have an effective width $W = H$. While this scheme will obviously have the advantages of simplicity the code itself will be much noisier. A larger effective width $W = 1.8 H$ is associated with shaped particles having the form of a Gaussian distribution (G.S.C.).

Hockney has shown that if an optimum time step Δ_t defined by

$$(\omega_p \Delta_t)_{opt} = \min(\frac{1}{2} H/\lambda_D, 1)$$

is chosen; then

$$\frac{\tau_H}{\tau_S}_{opt} = \frac{K_H}{(H/\lambda_D)^2} \quad (2)$$

where

$K_H = 5$ for NGP, 40 for CIC, 180 for G.S.C.

The noise in the system can be measured by the time average of the fluctuations in the electric field at a fixed point within the plasma. Hockney has shown that

$$\frac{\langle E_x^2 \rangle}{8\pi n m V_e^2} = \frac{0.12}{N(\lambda_D^2 + W^2)} \quad (3)$$

The most desirable simulation plasma is one in which τ_S is large, τ_H much larger and the noise level small. So we would, all things being equal, prefer a code with large heating ratio K_H , small noise level $\langle E_D^2 \rangle$ and containing particles of large effective width W .

8.3.2(c) Experimental Results

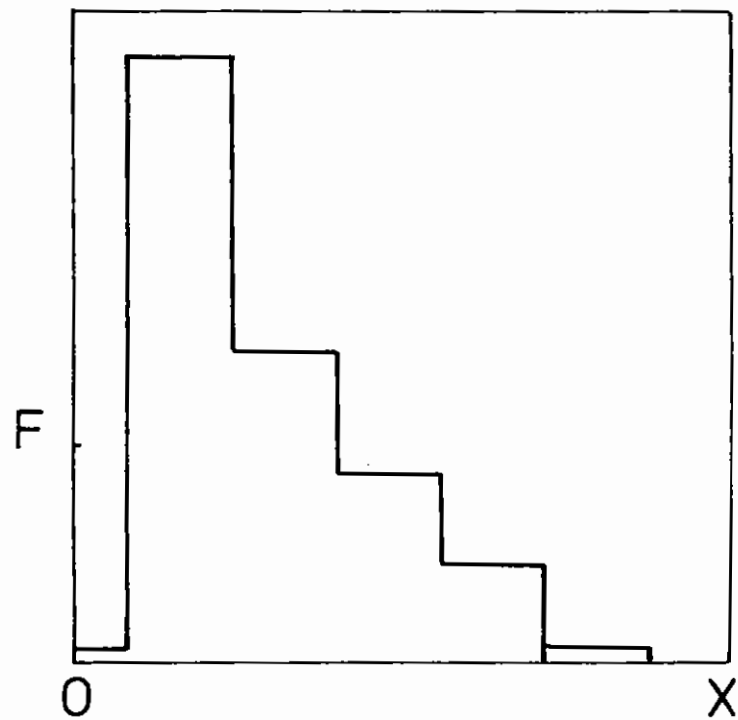
In all the simulations reported here the number of elements used was as follows: for the linear case $10 \times 10 = 100$ elements and for the cubic case $4 \times 4 = 16$ elements.

8.3.2(d) The law of force between particles

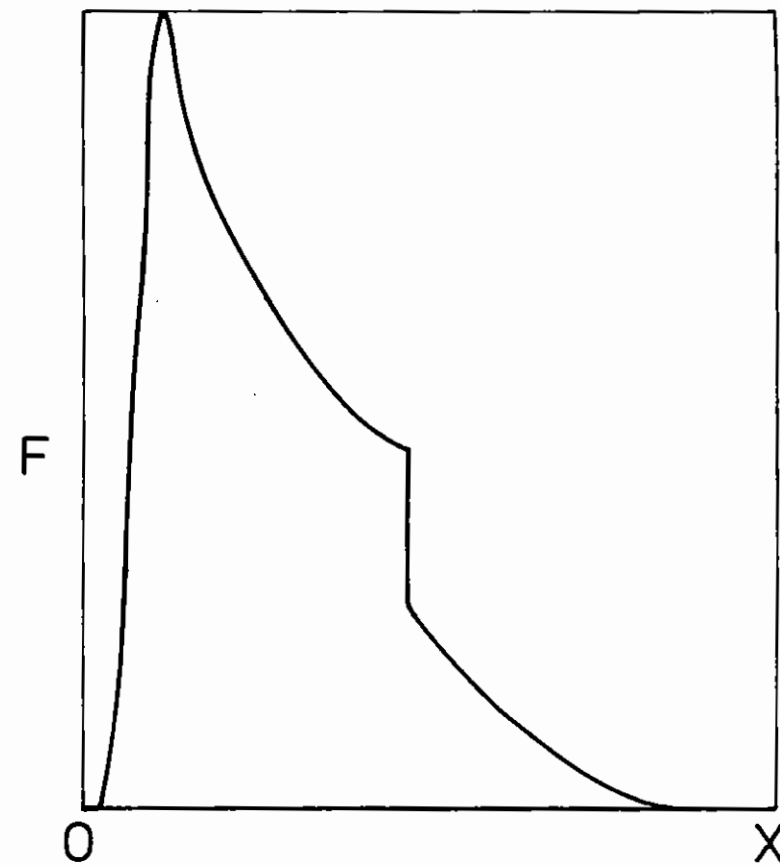
We have examined the law of force obtaining in our models. To do this one particle is placed, and held fixed, at a prescribed point. The vector $\underline{E}$ is calculated and the functional equation minimised to obtain nodal values of the potential ϕ . From these nodal values the field E_x is obtained. In Figures 8.11(a), (b) the force between two particles, one held at the origin, is plotted as a function of x . The force distribution in other directions differs only in detail from that given here. It is seen that on average the force falls off as $1/x$ as required but that there are abrupt changes in the magnitude as element boundaries are crossed. We also observe that the force falls to zero at the site of the fixed particle emphasizing that self forces do not exist in any of the models.

8.3.2(e) Collision and heating times

Figure 8.12 plots $\langle v_{||} \rangle$ against time, from which we may deduce the respective collision times. From these we see that for a simulation with $N_D = 16$ the collision times are $\sim 315 \omega_p^{-1}$ while for $N_D = 64$ they increase to $1200 \omega_p^{-1}$. The figures also show that collision times are unaffected by the finite element used. From (1) using these values for τ_S gives particle widths $W = 1.5 \lambda_D$ when $N_D = 16$, and $W = 1.4 \lambda_D$ when $N_D = 64$, for both the linear and cubic models.



(a) LINEAR ELEMENT



(b) CUBIC ELEMENT

Fig. 8.11 The force f (arbitrary units) between two particles. One particle is at the origin, the second is moved parallel to the x axis. Bars on the x axis denote element boundaries.

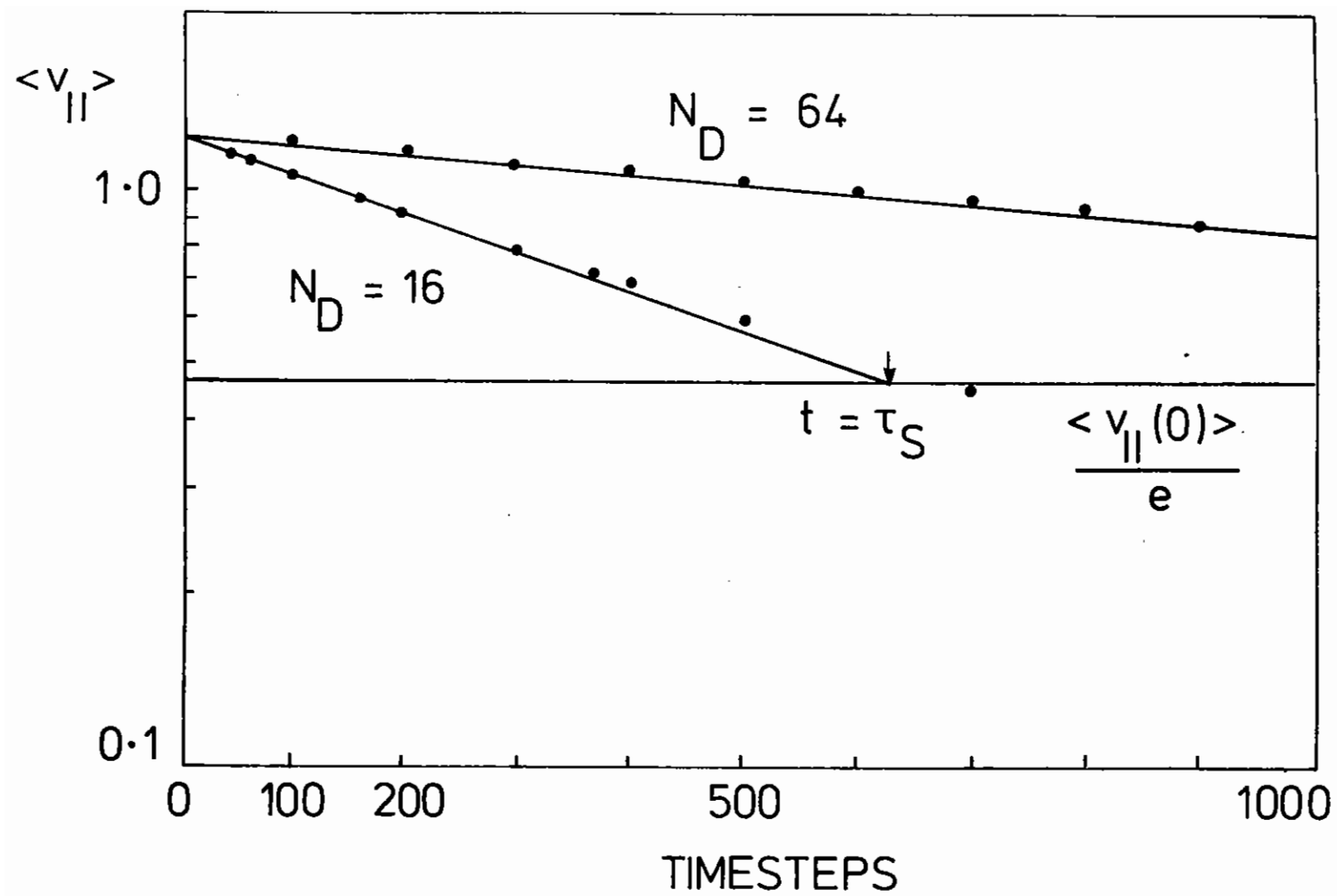


Fig. 8.12 Typical graph of $\langle v_{||} \rangle$ from which τ_S can be deduced. $\Delta t = 0.5 \omega_p^{-1}$

In Figure 8.13 $\langle h \rangle$ is plotted against time from which heating times τ_H can be deduced, but now the linear and cubic elements no longer lead to identical results. With a linear element and $N_D = 16$ $\tau_H = 1100 \omega_p^{-1}$ which increases to $\tau_H = 5500 \omega_p^{-1}$ when $N_D = 64$. The corresponding results for a cubic element are, $\tau_H = 2400 \omega_p^{-1}$ for $N_D = 16$, and, $\tau_H = 9500 \omega_p^{-1}$ for $N_D = 64$. Clearly using a cubic element is very advantageous in reducing the rate of energy increase in the simulation.

If we calculate τ_H/τ_S for the linear cases with $N_D = 16$, $N_D = 64$ we obtain $\tau_H/\tau_S = 8$. Identifying H in (2) with the element width, as seems reasonable, we estimate $K_H = 10$ (since $H = 1.6 \lambda_D$ for the linear element) whereas for the cubic element, $K_H = 125$. These results compare favourable with the values quoted earlier for finite difference codes.

Noise levels for linear elements satisfy (3), but the results obtained with cubic elements do not.

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8.3.3 The effect of magnetic field on Raman sidescatter

This section may be considered to be an extension of the work described in section 7.3.2 of the 1982 annual report. There, we considered Raman scatter (SRS) in which the emission occurred in the laser plane of polarization. However, the preferred polarization for SRS has the electric field of pump and scattered light aligned and, thus, it is more appropriate to consider emission out of the laser plane of polarization. Here we consider, in particular, exact sidescatter where the emission propagates parallel to a DC magnetic field, B_c (the laser light is assumed to be an extraordinary wave ie $k_0 \perp B_c$, $B_0 \parallel B_c$). The scattered light wave is then either a right or left circularly polarized electromagnetic wave. In the absence of the DC field, a linearly polarized pump wave generates a linearly polarized scattered wave which may be regarded as the sum of a left and right circularly polarized wave. With a finite DC field, the frequencies of the daughter waves and hence the locations of the decays become separate and distinct. Hence the DC

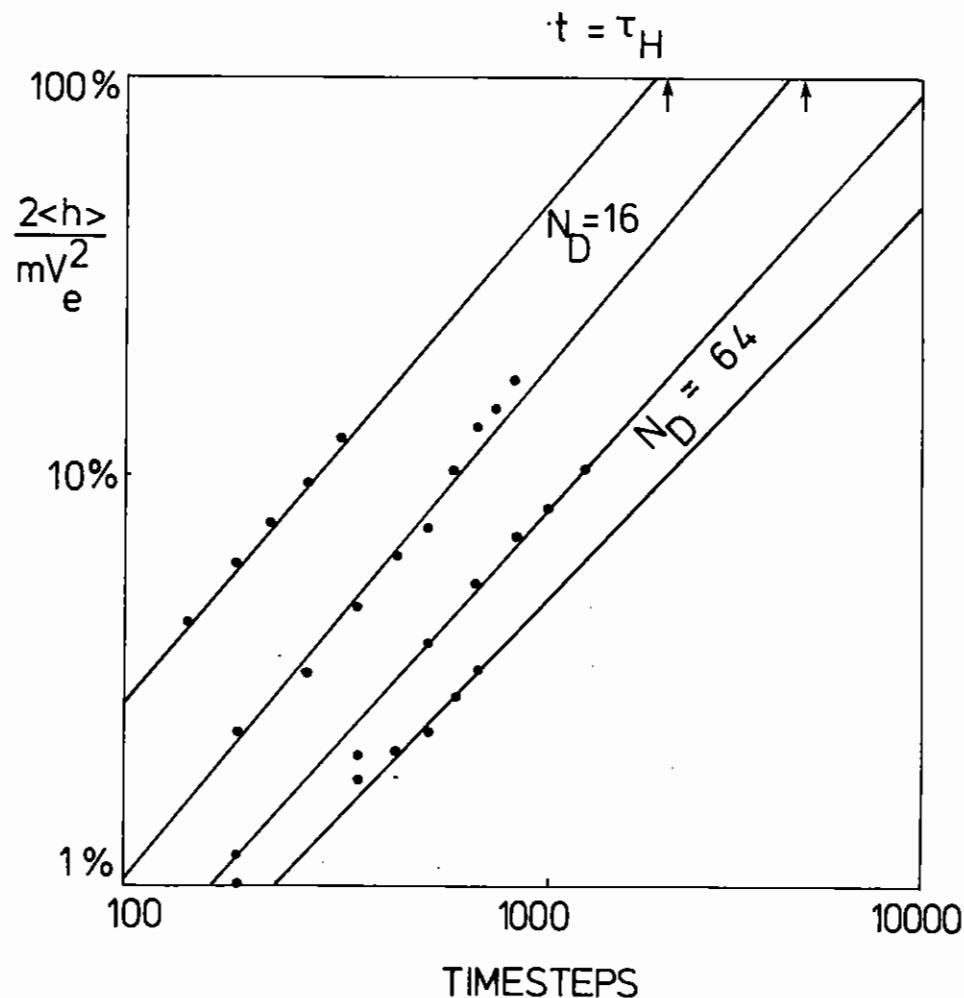


Fig. 8.13 Typical graphs from which the heating time τ_H can be deduced.

magnetic field gives rise to a splitting in the emission.

In stimulated Raman scattering the most unstable waves are generated near to the quarter-critical density of the plasma, where absolute instability takes place. Phase matching conditions demand that the scattered light-wave have a frequency slightly below $\omega_0/2$, shifted by an amount which depends on the plasma temperature:

$$\omega_{em} = \frac{\omega_0}{2} - \frac{9}{8} \frac{v_e^2}{c^2} \omega_0 \quad (1)$$

In the 1982 report it was shown that in the presence of a magnetic field this frequency becomes, for Raman backscatter,

$$\omega_{em} = \frac{\omega_0}{2} - \frac{9}{8} \frac{v_e^2}{c^2} \omega_0 \mp \frac{\Omega_e}{4} \quad (2)$$

and thus the magnetic shift acts to off-set the thermal shift. Similar frequency shifts arise when sidescattered light waves travel along magnetic field lines. The frequency and wave number matching conditions yield at the reflection point of the em waves:

$$\omega_{L,R} = \frac{\omega_0}{2} - \frac{9}{8} \frac{v_e^2}{c^2} \omega_0 \mp \frac{\Omega_e}{4} \quad (3)$$

In this expression $\omega_{L,R}$ represents the frequency of the left (L) or right (R) circularly polarized light-waves. Thus, a splitting of the $\omega_0/2$ light emanating from the plasma may occur, symmetrically spaced around the frequency point $\omega_0/2 - 9/8 v_e^2/c^2 \omega_0$. The separation between the red and blue wings of the spectrum is seen to be approximately $\Omega_e/2$ in frequency. The importance of this process to experiments is that it depends linearly on the magnitude of the magnetic field. In experiments performed at the Rutherford Appleton Laboratory by Sim and McGoldrick (8.27), and at Lawrence Livermore Laboratories by Phillion et al (8.28), such a splitting has in fact been observed in the $\omega_0/2$ light. In these

experiments the scattered light spectrum was mostly sampled along or close to the target normal. Paradoxically, this light can still be a product of Raman sidescatter taking place near $n_c/4$ as most of the light being generated there is collimated along the target normal due to refraction.

The above results can be used to deduce magnetic field strengths required to produce splitting of the $\omega_0/2$ light observed in experiments. For example, on aluminium targets irradiated by 0.35 μm laser light, Sim and McGoldrick (8.27) measured a separation of $\sim 80 \text{ \AA}$ between the red and blue wings. The magnetic field required in this case is $\sim 3.0 \text{ MG}$. Phillion et al (8.28) at Lawrence Livermore Laboratory observed a separation of $\sim 420 \text{ \AA}$ on 0.35 μm laser-irradiated gold targets. The magnetic field required to produce this separation is $\sim 7 \text{ MG}$.

A general fully electromagnetic dispersion relation has been derived and has been used to obtain the growth rates of the processes discussed above. The complexity introduced by the magnetic field has meant that we have had to obtain solutions numerically. In Figure 8.14 and Figure 8.15 we have plotted frequency and growth rates as functions of density (through ω_p). Figure 8.14 corresponds to the right-hand-circularly polarized waves and Figure 8.15 the left-hand-circularly polarized wave. The broken lines are for $\Omega_e/\omega_0 = 0.1$ and the solid lines for $\Omega_e/\omega_0 = 0.2$. An electron temperature $V_e/c = 0.1$ has been chosen.

Consider first of all Figure 8.14 for the RHCP waves. Note that the growth rates are larger at $\Omega_e/\omega_0 = 0.2$ than at $\Omega_e/\omega_0 = 0.1$. Note also that the frequencies of the scattered light waves never falls below $\omega_0/2$ and that the extension beyond $\omega_0/2$ is greatest when the magnetic field is largest. This behaviour is in agreement with equation (3) above. It can also be observed that the magnetic field pushes the instability towards lower densities (the cut-off density is lower at $\Omega_e = 0.2$ than at $\Omega_e = 0.1$). This can again be deduced using phase matching conditions and gives an expression which determines the maximum density at which decay into each of the two polarizations can occur:

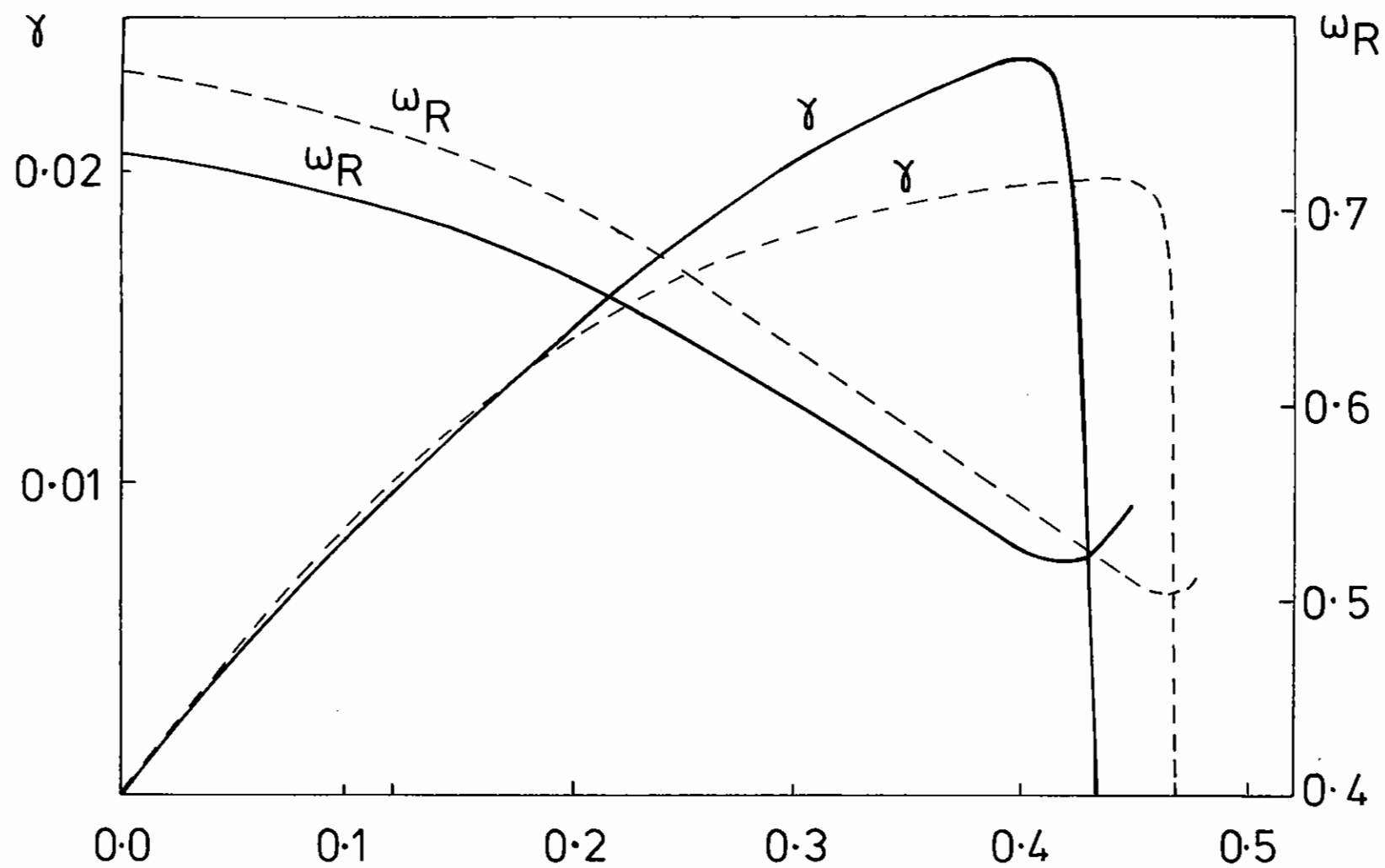


Fig. 8.14 Frequencies and maximum growth rates as a function of density (through ω_p) for stimulated Raman side scattering (RHCP).

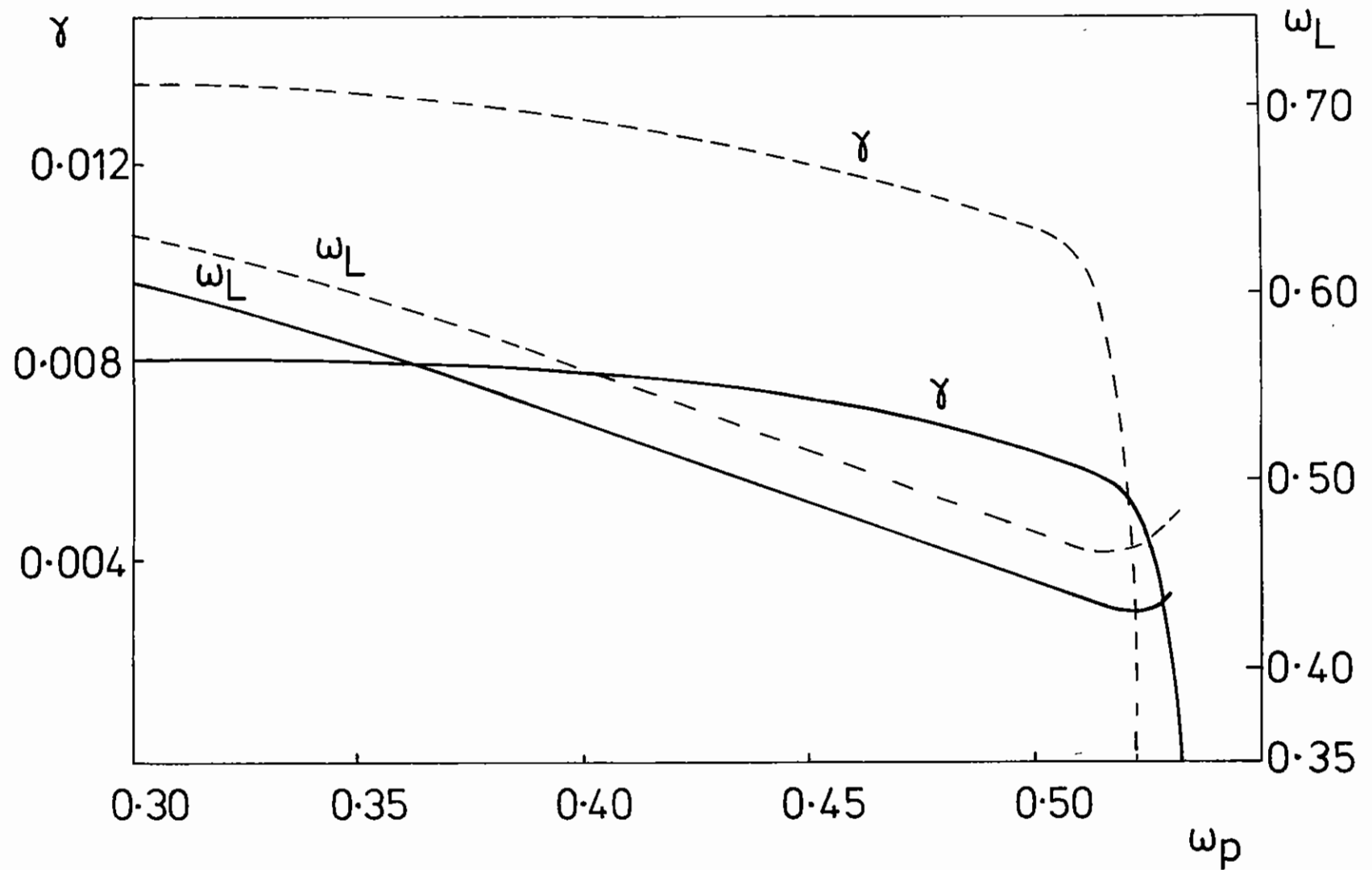


Fig. 8.15 Frequencies and maximum growth rates as a function of density (through ω_p) for stimulated Raman side scattering (LCHP)

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$$\omega_{\text{max}}^2 = \frac{\omega_0^2}{4} - \frac{9}{8} \frac{v_e^2}{c^2} \omega_0^2 + \frac{\Omega_e \omega_0}{4} \quad (4)$$

Consider now Figure 8.15 for the decay into the LHCP wave. It can be seen that the magnetic field inhibits growth of this wave, in contrast to the other polarization. This is not unexpected as in one case (RHCP) the electric vector rotates in the same direction as the electrons (about the magnetic field) while in the other (LHCP) it rotates in the opposite direction. There is therefore an "easy" direction (RHCP) and a "hard" direction (LHCP) for driving the electro-magnetic waves, so that the magnetic field produces opposite effects on each polarization. Note that for this polarization (LHCP) the instability can take place at densities lying above n_c/c ; its extension is proportional to the strength of the magnetic field, in agreement with equation (4). In this case the scattered light frequency lies below $\omega_0/2$ in the vicinity of $n_c/4$, as predicted by equation (3).

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8.3.4 Kinetic theory of Raman scatter in a magnetized plasma

In section 8.3.3 the effect of magnetic fields on stimulated Raman scattering was discussed. In this section we present results from a kinetic treatment. The Vlasov equation has been used to derive a nonlinear dispersion relation describing the decay of a pump extraordinary wave into an electrostatic Bernstein mode and a scattered extraordinary wave. The treatment is valid for arbitrary pump power, magnetic field strength and electron temperature.

The nonlinear dispersion relation has the form:

$$\begin{vmatrix} D_{xx} & D_{xy} & \beta_{xx} \\ D_{yx} & D_{yy} & \beta_{yx} \\ \alpha_{xx} & \alpha_{xy} & D_{xx} \end{vmatrix} = 0$$

The coupling terms α and β are products of the incident laser field in whose absence the dispersion relation can be written as:

$$D_{xx} (D_{xx} - D_{yy} - D_{xy} - D_{yx}) = 0$$

D_{xx} gives the dispersion relation of the Bernstein waves and $(D_{xx} - D_{yy} - D_{xy} - D_{yx})$ gives that of the scattered light waves. These terms may be found in chapter 8 of the standard text by Krall and Trivelpiece. The nonlinear dispersion relation is extremely complicated and it has been necessary for us to obtain solutions numerically. In general, the growth rates of the instabilities depend sensitively on both the strength of the magnetic field and Larmor radius of the electrons, as is expected. The Bernstein wave which starts on the branch of the dispersion relation corresponding to the upper-hybrid frequency $(\omega_p^2 + \Omega_e^2)$, behaves differently from other waves capable of being excited. In general the temperature and magnetic field effects can act to reduce growth of this wave. The inducement of Bernstein waves other than the upper-hybrid wave is enhanced by increasing the magnetic field strength and electron temperature.

In Figure 8.16 we have sketched the growth rate γ as a function of the wavenumber k of the Bernstein waves at a fixed density (given by ω_p), with frequencies normalized to ω_0^{-1} , wavenumbers to ω_0/c and speeds to $1/c$. In the figure the solid line represents $\Omega_e = 0.16$ (Ω_e is the cyclotron frequency) and the broken line $\Omega_e = 0.17$. The quiver velocity $V_0 = 0.1$ and the electron temperature $V_e = 0.1$. The numbers on the curves indicate the frequency the waves have at the position of maximum growth. At these field strengths only the first three Bernstein modes are capable of being excited ($n = 1, 2, 3$). The left-hand branches represent forward scattered light waves and the right-hand branches backscattered light waves. For $n = 3$ the forward and backward propagating waves have merged together, as they are close to their reflection point at this density. For $n = 1$ and $n = 2$ the forward and backward scattered waves are well separated as their point of reflection lies at higher ω_p . Note that growth of the backscattered waves is always larger than the forward scattered waves. In fluid treatments of Raman scattering γ is proportional to k . This is

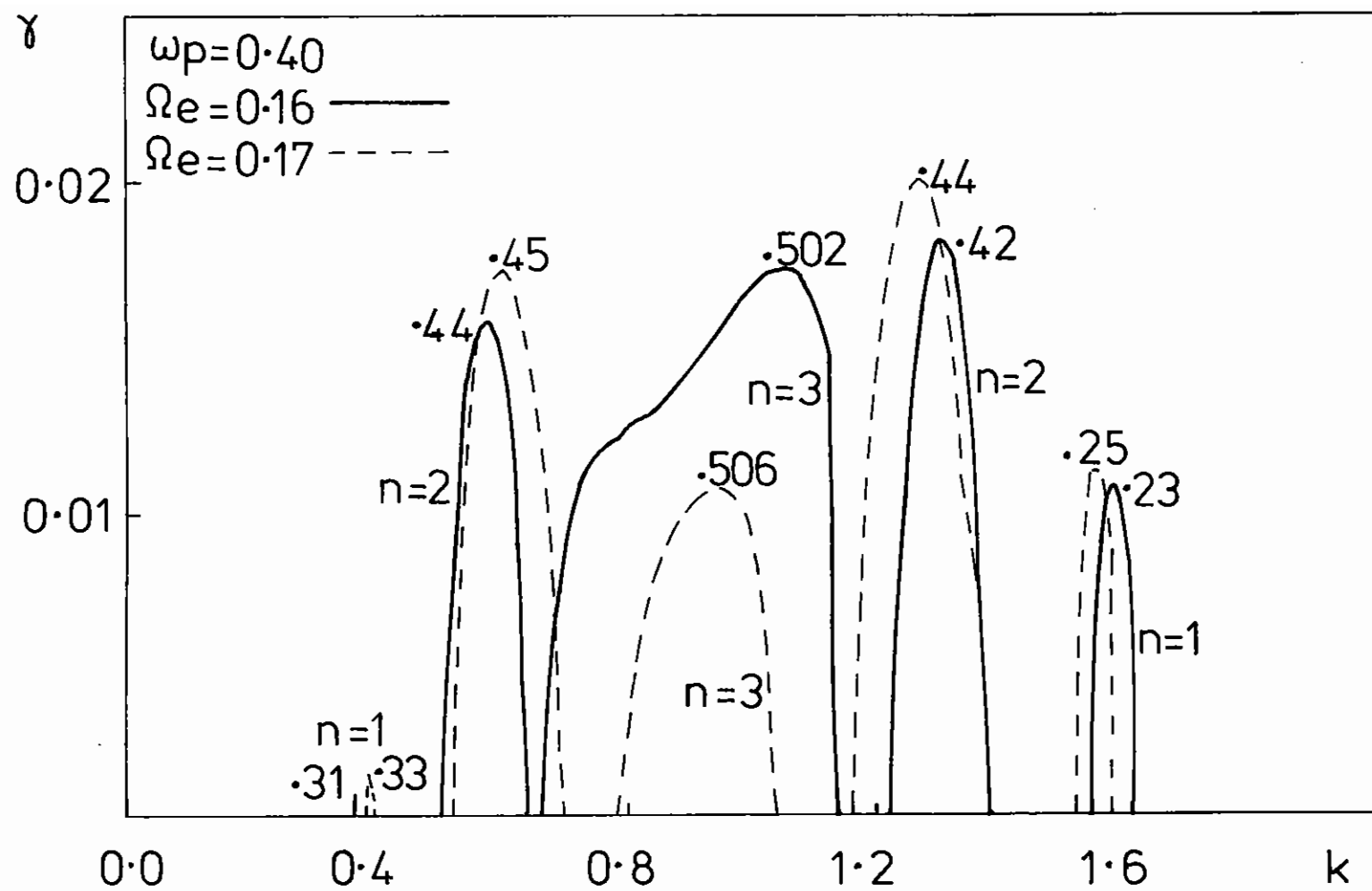


Fig. 8.16 Growth rates (γ) as a function of k for scattering off the first three Bernstein waves ($n = 1, 2, 3$).

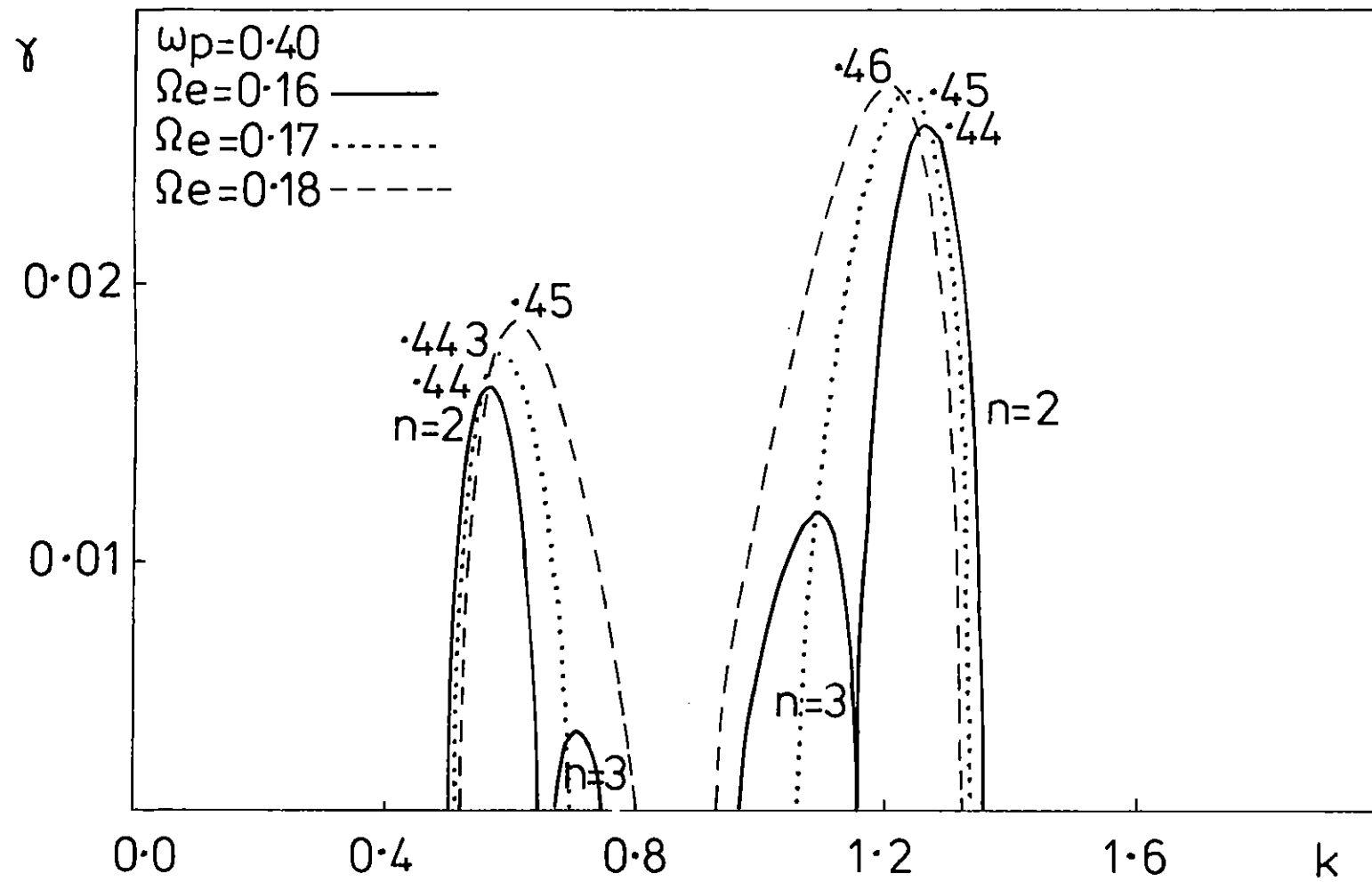


Fig. 8.17 Growth rates (γ) as a function of k for scattering off Bernstein waves ($n = 2, 3$).

evidently also true for scattering off Bernstein waves, as seen in the figure. As the magnetic field is increased each wave is pushed nearer to its reflection point, and thus the broken lines corresponding to a given branch (ie $n = 1, 2$ or 3) lie closer together than the solid lines. When $\Omega_e = 0.17$ the point of optimum phase matching has been passed for the $n = 3$ mode and consequently its growth rate is rapidly decreasing.

To investigate the effect electron temperature has on the instabilities consider Figure 8.17 where γ has been plotted against k for $\Omega_e = 0.16$ (solid lines), $\Omega_e = 0.17$ (dotted line), $\Omega_e = 0.18$ (broken line) and for $V_e = 0.06$. In this case the only branches of the dispersion relation which appear correspond to $n = 2$ and $n = 3$ (the growth rate of the $n = 1$ mode is $\sim$ zero). A large reduction in the growth rate of the $n = 3$ mode is seen to have occurred as a consequence of the decrease in the magnitude of the Larmor radius of the electrons. Contrasting this is the increase in the growth rate of the $n = 2$ mode. The frequency characteristics of this branch coincide with the upper-hybrid wave and thus it may be considered to be the usual fluid mode including kinetic effects.

Of particular interest to experimentalists are the frequencies of the scattered light waves associated with these instabilities. A red shift in $\omega_0/2$ light is known to be a signature of Raman scattering. This normally only takes place near the quarter-critical density of the plasma. In Figure 8.16 we observe that even for $n = 16n_{cr}$ ($\omega_p = 0.40$) Bernstein waves can give rise to a red-shifted spectrum, eg, when $\Omega_e = 0.16$ the $n = 3$ mode frequency is approximately equal to 0.502, implying that the associated light frequency be ~ 0.498 . In general the frequencies depend on V_e , Ω_e and ω_p .

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8.3.5 The effect of Brillouin scatter on Raman scatter

Most experiments in which stimulated Raman scattering (SRS) occurs will also observe stimulated Brillouin scattering (SBS). Therefore, the high

frequency waves of SRS propagate against a background of finite amplitude ion waves.

We have studied the effect of a finite amplitude ion wave on SRS primarily at low densities $\omega_p \ll \omega_0$ appropriate to an experiment at UCLA (8.29). It can give rise to (i) significant frequency shifts (ii) reduced growth and (iii) new instabilities having growth rates comparable with SRS.

Consider, for simplicity, 1-D Raman backscatter occurring against a zero-order density profile

$$n_0(x) = N_0(1 + \epsilon \cos 2k_0 x) \quad (1)$$

This represents the ion wave, wavenumber $2k_0$, amplitude ϵN_0 generated by SBS. A small static field is neglected, valid when $2k_0 \lambda_D < 1$. The ripple predominantly affects the electron plasma wave; its effect on light waves is proportionately $(V_e/c)^2$ less and is therefore neglected.

To elucidate the new physics introduced by the ripple, consider the plasma wave propagation described by

$$\left[\frac{\partial^2}{\partial t^2} + \omega_p^2(x) - 3V_e^2 \frac{\partial^2}{\partial x^2} \right] E(x, t) = 0 \quad (2)$$

where $\omega_p^2(x) = \omega_p^2(1 + \epsilon \cos 2k_0 x)$. Assuming $E \propto \exp(i\omega t)$ and letting $z = k_0 x$, $a = (\omega_p^2 - \omega^2)/3k_0^2 V_e^2$, $q = \epsilon/6k_0^2 \lambda_D^2$ we obtain Mathieu's equation

$$\frac{d^2 E}{dz^2} + (a - 2q \cos 2z) E = 0 \quad (3)$$

Solutions naturally divide themselves into those for small q and those for large q . Small q means that the variation in $\omega_p^2(x)$ due to the ripple can be compensated for by a change in k in the dispersion relation $\omega^2 = \omega_p^2(x) + 3k^2 V_e^2$. Consequently the plasma wave is no longer sinusoidal but still propagates at all points. Large q means that the pressure term $3k^2 V_e^2$ is too small and cannot compensate the variation in $\omega_p^2(x)$. Plasma

waves then become trapped in the troughs of the density ripple and are evanescent through the ripple peaks.

An infinite set of normal modes exist comprised of different linear combinations of Fourier modes $k \pm 2nk_0$; for, Fourier analyzing (3) yields

$$(a - K^2) E(K) = q[E(K + 2) + E(K - 2)] \quad (4)$$

where $K = k/k_0$. q is the natural coupling parameter and is largest in cool, dense, highly rippled systems. The coupling into high k , low phase velocity modes will result in increased Landau damping. Also, even if $2k_0\lambda_D < 1$, $2nk_0\lambda_D > 1$ for some n and a kinetic treatment is needed.

The laser coupling introduces the possibility of new instabilities which may be understood as follows. Suppose the light waves (ω_0, k_0) and $(\omega_0 - \omega, k_0 - k)$ drive a disturbance (ω, k) . This disturbance is coupled by the ripple to disturbances at $(\omega, k \pm 2nk_0)$. In the absence of the ripple $k \pm 2nk_0$ correspond to an infinite set of plasma waves only one of which (ω, k) can phase match and be driven unstable by the light waves. With a finite amplitude ripple, if we frequency match with any of the normal modes they possibly may be driven unstable since they are all now coupled to the laser. In practice, however, only a few are unstable.

The kinetic analogue of (4), including the laser coupling, is

$$\left[1 + \frac{1}{\chi(k, \omega)} - \frac{k^2 v_0^2}{D(k, \omega)}\right] E(k, \omega) = -\frac{\varepsilon}{2} [E(k + 2k_0, \omega) - E(k - 2k_0, \omega)] \quad (5)$$

where $(\omega, k) \equiv (\omega - \omega_0, k - k_0)$, $D(k, \omega) = k^2 c^2 + \omega_p^2 - \omega^2$ and χ is the usual electron susceptibility which for a Maxwellian plasma is

$$\chi(k, \omega) = -\frac{1}{2k^2 \lambda_D^2} Z'(\zeta) \quad (6)$$

$Z(\zeta)$ is the Plasma Dispersion Function and $\zeta = \omega/2k v_e$. When $\varepsilon = 0$, (5) gives the familiar SRS dispersion relation for 1-D scattering.

The solutions of (5) for the three most unstable roots are plotted in Figures 8.18 and 8.19. When $q = 0$, the mode ω_1 is the usual SRS plasma wave with $k \approx 2k_0$ while ω_2 has $k \approx -2k_0$ and ω_3 $k \approx 0$. As q increases and energy is channeled into the sidebands, so the growth of ω_1 is reduced while growth of ω_2, ω_3 becomes possible. For large q the growth of all three modes is limited by the increased Landau damping into the low phase velocity modes. As can be seen, ω_3 eventually dominates since it contains as its major component the high phase velocity mode $k \approx 0$. The separation in frequency increases with q .

Although we have not proved it, it seems intuitively reasonable to expect that as q increases and the plasma wave energy becomes more localized to, trapped in, the ripple troughs that the SRS convective instability may become absolute.

The implications of this analysis for high densities and for parameters more appropriate to laser-target experiments is currently under investigation.

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8.4 Transport Processes

8.4.1 Grad's 13 moment equations and electron thermal transport

The breakdown of the classical Spitzer-Härm theory of heat conduction (8.30) in laser-produced plasmas has resulted in the introduction of an ad-hoc flux limit in most hydrodynamic codes. A less arbitrary resolution of the problem of heat transport may be found by resort to the 13-moment equations derived by Grad (8.31).

By expanding the distribution function about a local Maxwellian in a series of Hermite polynomials in $\underline{v}$ and taking velocity moments of the Boltzmann equation Grad arrived at an alternative set of fluid equations. The important distinction between this approach and Chapman-Enskog theory is that the expansion used does not depend directly on the electron mean free path being much smaller than the scale length of variation of the

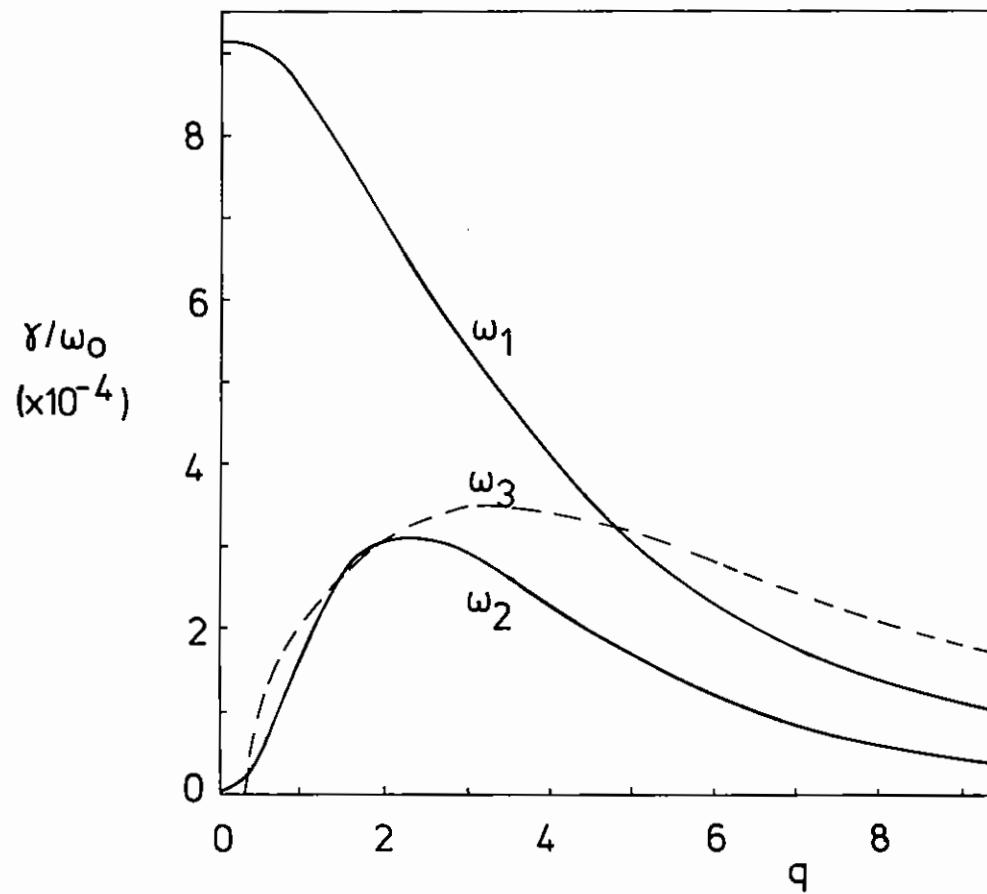


Fig. 8.18 Growth rate versus q for $v_0/c = 0.003$, $v_e/c = 0.01$, $n/n_c = 0.01$.

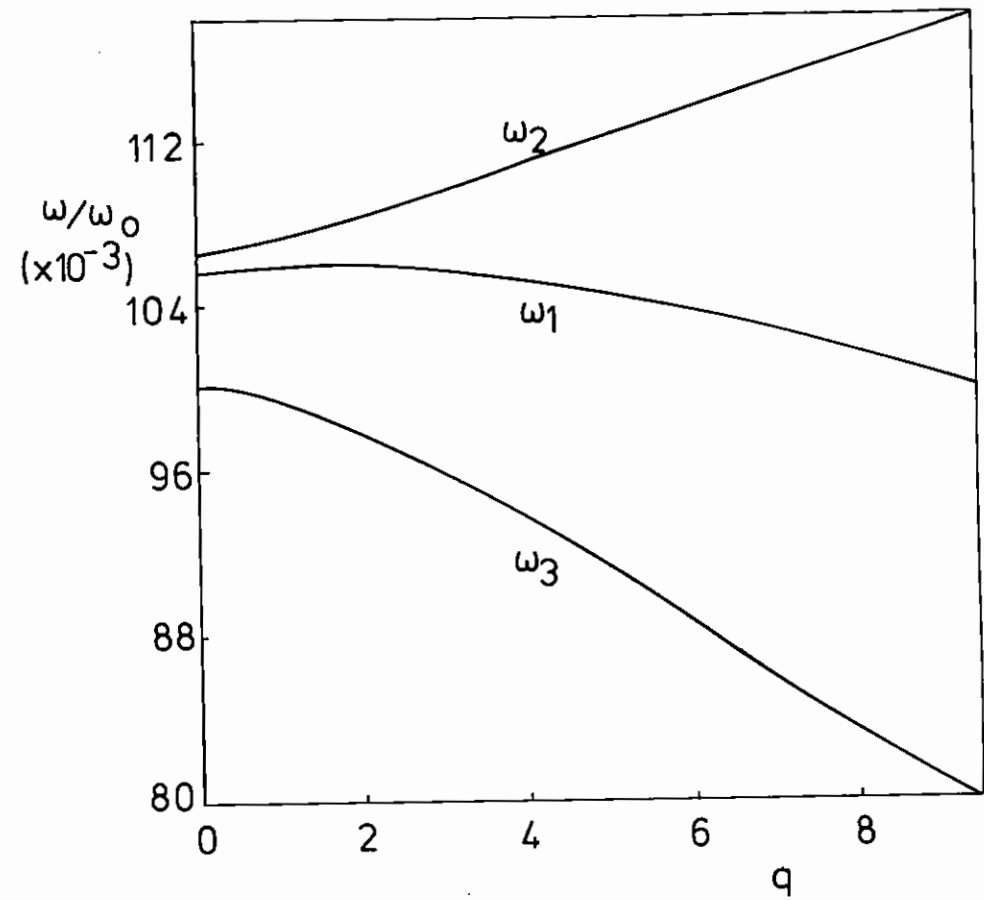


Fig. 8.19 Frequency versus q for $v_0/c = 0.003$, $v_e/c = 0.01$, $n/n_c = 0.01$.

macroscopic properties of the fluid. Thus in the regime of low collision frequencies and steep temperature gradients we can hope for a more realistic insight into transport physics from these equations.

In 3 dimensions there are 13 equations for 13 velocity moments of the distribution function, but our model is restricted to one dimensional slab geometry for which there are 5 independent variables. The first three moments are simply the density n , flow velocity u and pressure of the electron fluid p ; π and q , the viscosity and heat flux respectively, are defined by:

$$\begin{aligned}\pi &= mn \langle (v_x - u)^2 \rangle - p \\ q &= \frac{mn}{2} \langle (v_x - u)(v_x - u)^2 \rangle\end{aligned}\quad (1)$$

The set of equations to be solved is:

$$\begin{aligned}\frac{\partial n}{\partial t} + \frac{\partial}{\partial x} (nu) &= 0 \\ \frac{\partial u}{\partial t} + u \frac{\partial u}{\partial x} + \frac{1}{mn} \frac{\partial}{\partial x} (p + \pi) + \frac{e}{m} E &= \left(\frac{\partial u}{\partial t} \right)_c \\ \frac{\partial p}{\partial t} + u \frac{\partial p}{\partial x} + \frac{5}{3} p \frac{\partial u}{\partial x} + \frac{2}{3} \pi \frac{\partial u}{\partial x} + \frac{2}{3} \frac{\partial q}{\partial x} &= \left(\frac{\partial p}{\partial t} \right)_c \\ \frac{\partial \pi}{\partial t} + u \frac{\partial \pi}{\partial x} + \frac{7}{3} \pi \frac{\partial u}{\partial x} + \frac{4}{3} p \frac{\partial u}{\partial x} + \frac{8}{15} \frac{\partial q}{\partial x} &= \left(\frac{\partial \pi}{\partial t} \right)_c \\ \frac{\partial q}{\partial t} + u \frac{\partial q}{\partial x} + \frac{16}{5} q \frac{\partial u}{\partial x} + \frac{(5p + 7\pi)}{2m} \frac{\partial}{\partial x} \left(\frac{p}{n} \right) \\ + \frac{p}{mn} \frac{\partial \pi}{\partial x} - \frac{\pi}{mn} \frac{\partial}{\partial x} (p + \pi) &= \left(\frac{\partial q}{\partial t} \right)_c\end{aligned}\quad (2)$$

The collisional contributions on the right hand sides are evaluated in terms of the 5 moments and a "collision frequency" ν , according to the model used. Salat (8.32) has given these terms for a Fokker-Planck

collision model.

Since we are hoping to use these equations in conditions under which collisional effects are small we have tested the model in a collisionless regime. Grad's expansion should be valid over timescales smaller than a collision time provided the distribution function remains "near-Maxwellian".

Assuming quasi-neutrality and no collisions the set of equations (2) reduces to 3 equations describing heat conduction in the electron fluid:

$$\begin{aligned}\frac{\partial p}{\partial t} + \frac{2}{3} \frac{\partial q}{\partial x} &= S \\ \frac{\partial \pi}{\partial t} + \frac{8}{15} \frac{\partial q}{\partial x} &= 0 \\ \frac{\partial q}{\partial t} + \frac{5}{2} \left(\frac{p+\pi}{mn} \right) \frac{\partial p}{\partial x} + \left(\frac{p-\pi}{mn} \right) \frac{\partial \pi}{\partial x} &= 0\end{aligned}\quad (3)$$

S is a source term for the electron pressure. A slab of plasma of length 3 microns, density 10^{21} cm^{-3} and initial temperature 1 keV was heated at one end using a source term

$$\begin{aligned}S &= \frac{P_0}{\tau} (1 - X) \quad X < 1 \\ &= 0 \quad X > 1\end{aligned}\quad (4)$$

in which P_0 is the initial electron pressure, τ is the time simulated (0.5 ps) and X is measured in microns. Equations (3) were solved numerically with boundary conditions $q = 0$ at both ends. An identical experiment was performed with a $1\frac{1}{2}$ dimensional electrostatic particle-in-cell code, in which the boundaries were reflective and the electrons were heated in such a way as to keep their distribution Maxwellian in the absence of heat flow.

Figure 8.20 shows the spatial variation of the three variables p/p_0 , π/p and q/q_{fs} where q_{fs} is the free streaming limit.

The solid line represents the 13 moment solution and the circles give the results from the particle code. The agreement between the two for the p and q plots is satisfactory and what differences there are could well have been predicted. In the fluid model a thermal disturbance can propagate no faster than the thermal speed in the unheated region. By contrast, in the particle simulation the heat flux profile has a lower maximum and has spread further.

For the variable π the codes are not in good agreement but since π/p is small the effect on the p and q results is not significant. This discrepancy is to be attributed to the 13 moment equations where the expansion of the distribution function is truncated with only 3 of the third order velocity moments included.

The flux of p is accurately treated as $\frac{2}{3}q$, but the flux of π is not an independent variable in Grad's equations and is approximated to $\frac{8}{15}q$. This is generally an overestimate. The lesson to be drawn from this is to be circumspect about results obtained when π/p is large.

In a collision-dominated plasma the 13 moment equations (2) were applied to both ions and electrons in a one fluid, two temperature Lagrangian code.

We used Fokker-Planck collision terms as derived by Salat (3), which for the electrons are:

$$\left(\frac{\partial p_e}{\partial t}\right)_c = 2\nu_{ei} \frac{m_e}{m_i} n_e (T_i - T_e)$$

$$\left(\frac{\partial \pi_e}{\partial t}\right)_c = -\frac{6}{5} \nu_{ei} \pi_e$$

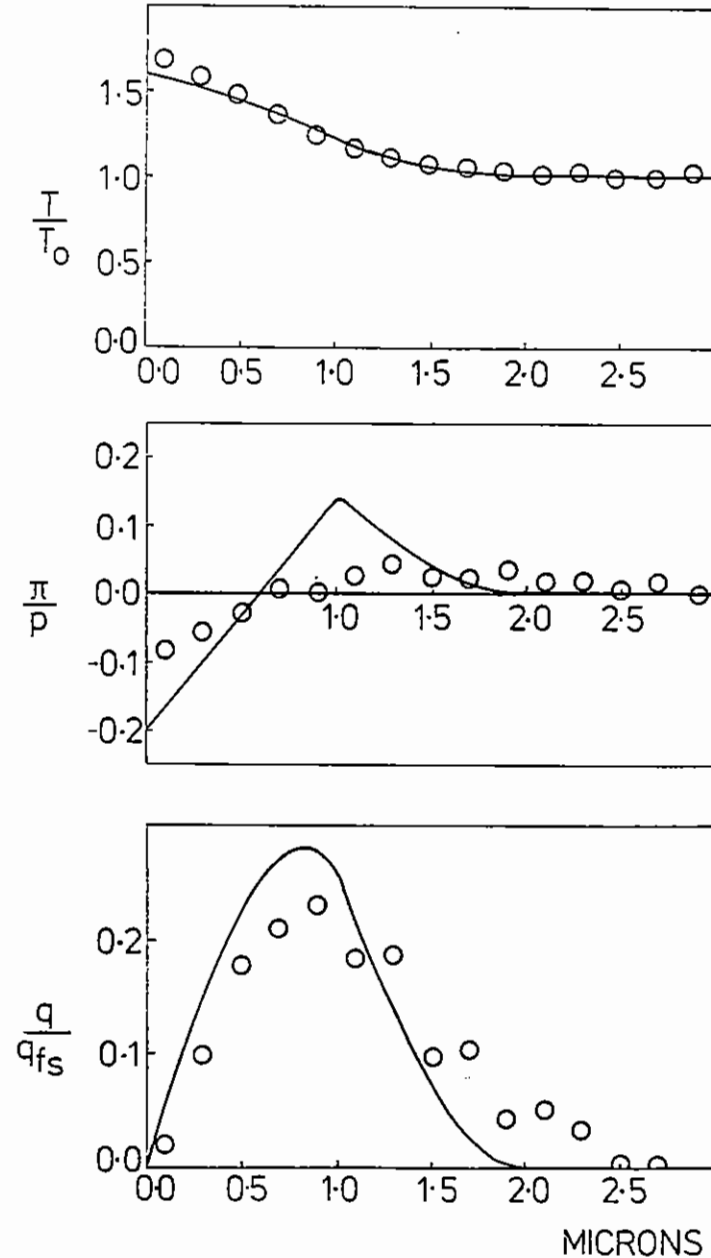


Fig. 8.20 Plots of p/p_0 , π/p , q/q_{fs} against x at the end of the simulation. The full lines give the results of the 13 moment equations, and the circles give the results for the particle code.

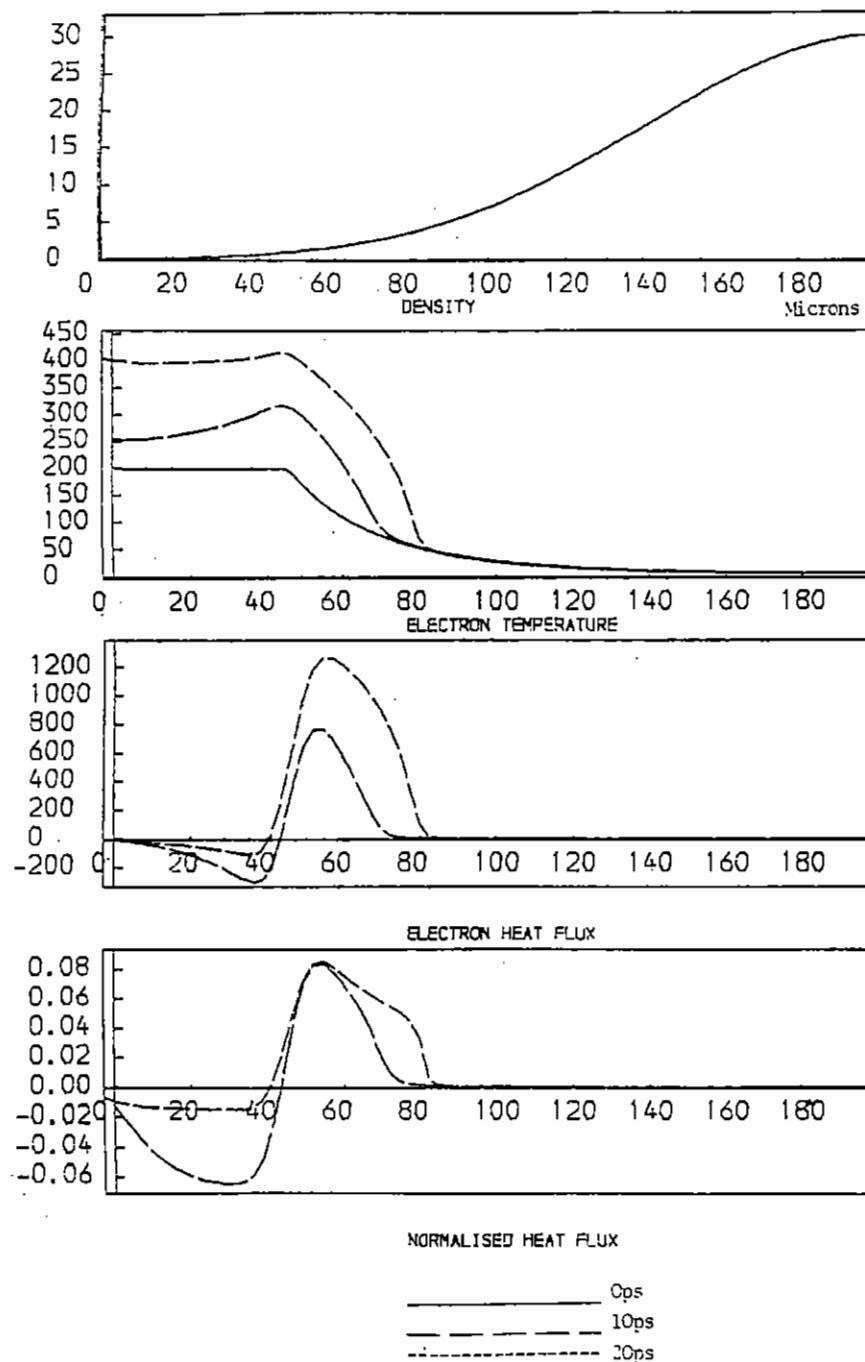


Fig. 8.21 Spatial profiles at various times throughout the simulation.

$$\left(\frac{\partial q_e}{\partial t}\right)_c = -\frac{13}{10} \nu_{ei} q_e$$

with the electron-ion collision frequency given by

$$\nu_{ei} = \frac{8}{3} \left(\frac{\pi}{2}\right)^{\frac{1}{2}} \frac{n_i Z^2 e^4}{m_e^2 T_e^{3/2}} \log \Lambda \quad (5)$$

where $\log \Lambda$ is the Coulomb logarithm.

Results were obtained from a run in which an initial density profile of the form shown in Figure 8.21 was used. $q = 0$ is imposed at both boundaries. Plots of n_e , T_e , q_e and q_e/q_{fs} as functions of x are shown in Figure 8.21 for times 0, 10 and 20 ps into the run. These variables are measured in code limits as follows:

$$\begin{aligned} n_e &= 10^{21} \text{ cm}^{-3} \\ T_e &= 5.69 \text{ eV} \\ q_e &= 0.91 \times 10^{18} \text{ gs}^{-3} \end{aligned} \quad (6)$$

Experiments using this 13 moment code over a range of initial conditions are continuing.

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8.4.2 Nonlinear heat flow

8.4.2(a) Electron heat transport in a steadily ablating Lorentz plasma

The breakdown of Spitzer's local perturbation theory of heat flow in laser-plasmas (8.33) has motivated numerical studies of electron heat transport at the kinetic level (8.12) (8.34) (8.35). The kinetic approach has the virtue of being able to describe non-local effects and has provided useful insights into electron transport.

Common to all these studies (8.12) (8.34) however, is the assumption of a stationary plasma (ie zero ion motion). We describe here a Fokker-Planck study of electron transport with self-consistent ion motion. We consider a slab of a steadily ablating plasma. The temperatures of the boundaries are fixed and the plasma is accelerated down a pressure gradient from the 'cold' to the 'hot' boundary.

A multi-(total energy) groups discrete ordinate code is used to obtain the simultaneous solutions of the steady state 1-D electron Fokker-Planck equation and the plasma equation of motion, for a given temperature ratio at the boundaries and mass flux through the system. The plasma carries no electric current. As an illustration of the model Figure 8.22 shows the kind of macroscopic profiles generated by the code. Results are still in a very preliminary stage.

One interesting point to emerge is the non-existence of a monotonic electric potential solution for high (but still subsonic) plasma flow velocities. From the ion equation of motion

$$Zn_1 eE = m_1 n_1 v_1 \frac{\partial v_1}{\partial x} + \frac{\partial P_1}{\partial x} + R_{1e}$$

(R_{1e} is the rate of change of momentum due to collisions with electrons) one sees that if there were no ion motion the electric field is simply the thermoelectric field which is directed down the temperature gradient. If however the convection term was dominant, the electric field would change sign ($v_1 \frac{\partial v_1}{\partial x} / \frac{\partial T}{\partial x} > 0$, refer Figure 8.22). The plasma flow velocity is strongest near the 'hot' or 'critical' surface and here the electric field can change sign acting to accelerate the 'hot' electrons into and the ions out of the ablation region. The existence of a potential 'hump' can also act to keep some electrons trapped just inside the critical surface.

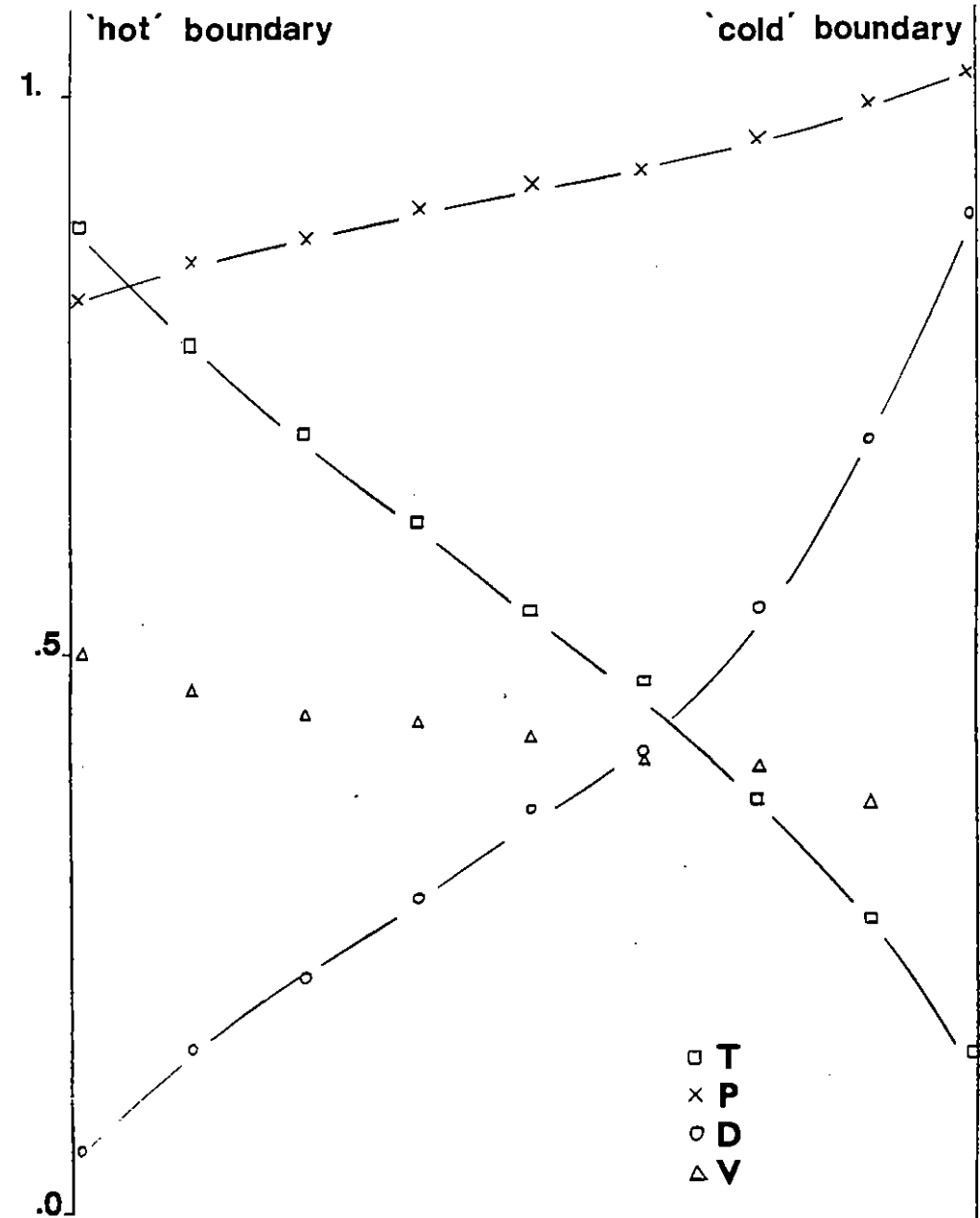


Fig. 8.22 Temperature (T), pressure (P) and density (D) profiles generated by the code for a boundary temperature ratio of 2. The units are arbitrary. The plasma flow speed (V) is normalised by the local sound speed. The direction of flow is from the 'cold' to the 'hot' boundary.

8.4.2(b) Transport and higher order Legendre moments of the electron distribution function

One method of solving the electron Fokker-Planck equation involves expanding the distribution function in spherical harmonics (or Legendre polynomials in 1-D)

$$f(x, v, \mu) = \sum_{\ell=0}^{\infty} P_{\ell}(\mu) f_{\ell}^{\ell}(x, v)$$

This produces an infinite set of coupled equations for the f_{ℓ}^{ℓ} 's. In practice, the expansion must be terminated at some point and to expedite computation it is desirable to keep as few terms as are deemed necessary.

It is seldom certain a priori how much influence the neglected higher Legendre moments have on transport. In a recent study with a 1-D polynomial expansion code Matte and Virmont (8.36) claimed that the first two terms (f^0 , f^1) are sufficient for electron transport. They found, using a 'hot' and 'cold' boundaries model, that the steady state heat flow differs from the Spitzer values because of the deviation of f^0 from a Maxwellian.

We look at these findings using the same model for the special case of a stationary, Lorentz plasma with no electric field or net current but with an appropriate self-consistent density gradient. The deviation of f^0 from a Maxwellian in a model of this type arises in part from boundary conditions. A boundary temperature ratio of 2 and above was used in (8.36). We try to minimise this effect here by using a ratio of 1.1. In the collisionless limit f^0 deviates very little from a Maxwellian.

The Fokker-Planck equation is solved using a multi-group discrete ordinate code. 50 angular direction ($0 < \mu_1 < 1$) and 12 energy groups at 0.5, 1.5, 2.5, ..., 11.5 T_c (in units of the temperature at the 'cold' boundary) are used. The electron mean-free-path is velocity and density dependent ($\propto v^4/n_1(x)$).

Figure 8.23 shows the steady state heat flow is significantly below the Spitzer value even for quite modest values of λ/L_T (~ 0.01). λ is the mean-free-path of an electron of energy kT and L_T is the temperature scalelength. (All quantities refer to the local values unless otherwise specified.)

Using a point 'A' on Figure 8.23 as an illustration, table 8.03 shows the values of f^1/f^0 and f^2/f^0 calculated from the distribution function obtained from the code. In column four (under 'A') is shown the values of $(\log f^0(E) - \log f^0(E + \Delta E))/\Delta E$ ie the gradient of $\log f^0$ v.s. $E (= \frac{1}{2}mv^2)$. This is a constant for a Maxwellian. One sees that $|f^1/f^0| < 1$, $|f^2/f^0| \ll 1$ and f^0 is very nearly a Maxwellian. These conditions are typical of many of the simulations for different collisionalities. In particular the less collisional the system, the more f^0 approximates a Maxwellian (because of the very small difference in boundary temperatures). One might have expected that under these conditions the Spitzer formulation would be a good approximation. From Figure 8.23 we see this is not the case.

Table 8.04 compares the flux-carrying part of the distribution f^1 calculated in the following ways:

- (a) f_{code}^1 , from the code,
- (b) $f_{sp}^1 = -\frac{\lambda}{2} \left(\frac{v^2}{v_c^2} - 4 \right) \frac{1}{T} \frac{\partial T}{\partial x} f_M$; ie the 'Spitzer f^1 '. f_M is a Maxwellian at the local density and temperature.

From the exact equation for f^1 ,

$$f^1 = -\frac{\lambda}{2} \frac{\partial f^0}{\partial x} - \frac{\lambda}{5} \frac{\partial f^2}{\partial x} \quad (1)$$

we define

$$f_0^1 = -\frac{\lambda}{2} \frac{\partial f^0}{\partial x} \text{ code} \quad (1a)$$

$$f_{02}^1 = -\frac{\lambda}{2} \frac{\partial f^0}{\partial x} \text{ code} - \frac{\lambda}{5} \frac{\partial f^2}{\partial x} \text{ code} \quad (1b)$$

In equations (1a), (1b) f_{code}^0 and f_{code}^2 are obtained from the code.

Table 8.04 shows that despite the very small ratio of f^2/f^0 here, neglecting the f^2 term in equation (1) can make a factor of 2 changes in the results. The heat flow here, calculated with f_0^1 is twice the actual result. For more collisional systems the contribution from the f^2 term decreases but the ratio of f^2/f^0 remains about the same. In this case at least, the relative sizes of f^0 and f^2 do not correlate with the influence of f^2 on transport.

f_{sp}^1 overestimates the flux by up to an order of magnitude for the more energetic electrons. Replacing f_{code}^0 in equation (1a) with a Maxwellian at the local temperature and density gives essentially the same results at f_{sp}^1 .

The reason f_{sp}^1 overestimates the heat flow is that the 'effective' temperature and density gradients of the energetic heat carrying electrons can be much smaller than the gradients for these quantities which are obtained by integrating over the whole energy spectrum. The spatial variation in the local temperature and density are caused mainly by the spatial variation in the number density of relatively slow collisional electrons which carry little flux. The energetic, flux-carrying electrons are essentially streaming through the system (in both directions) suffering relatively few collisions because of the v^4 dependence of the mean-free-path. Replacing f^0 by some (isotropic) distribution (eg a Maxwellian) which is governed by the local density and temperature will therefore give a wrong estimate of the anisotropy of these energetic electrons. This agrees with the study by Matte and Virmont (8.36).

However, we note that it is not the deviation of f^0 from a Maxwellian that is in itself responsible for the difference between the actual and the Spitzer heat flows; this can be seen, for example, by comparing the

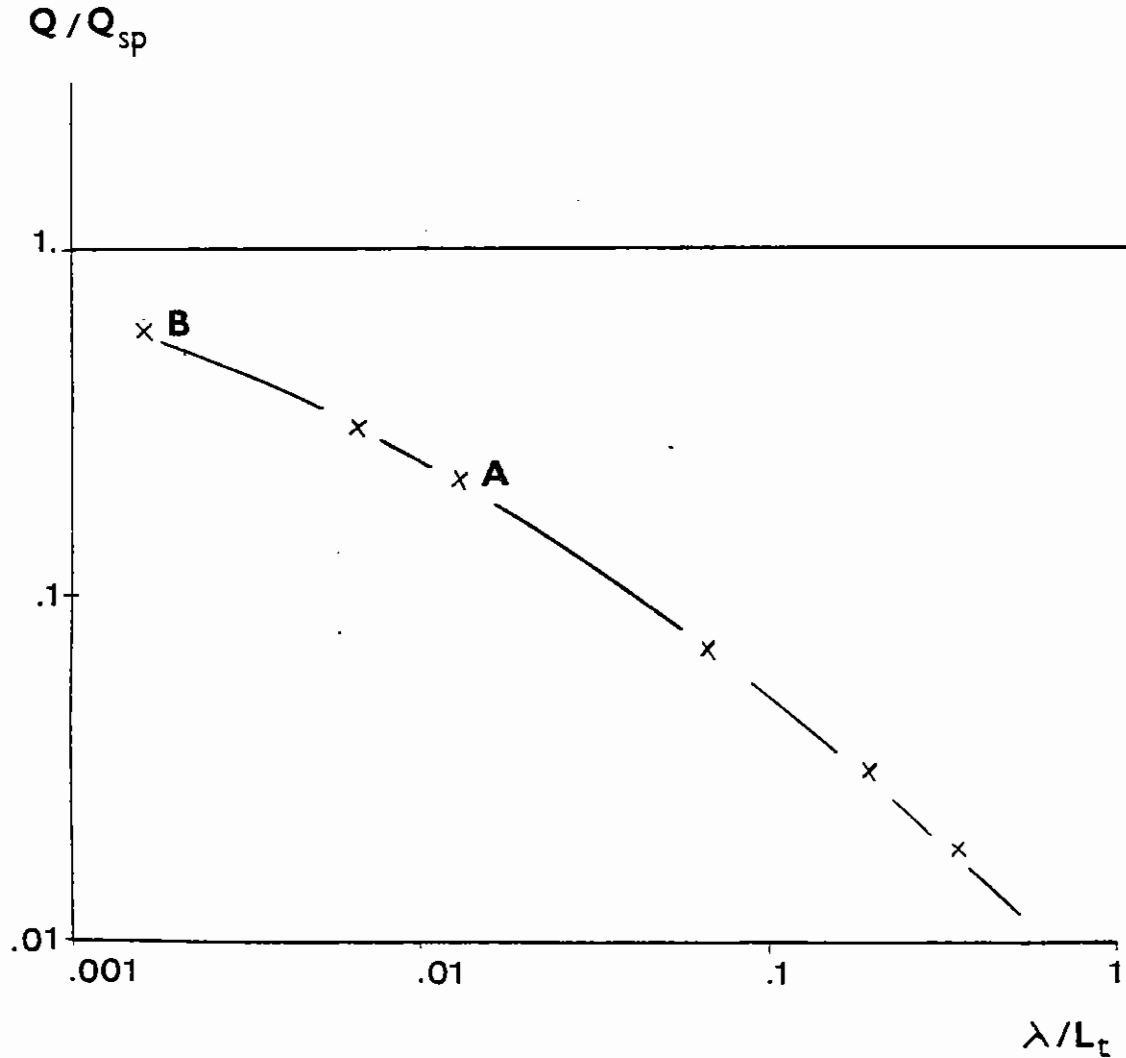


Fig. 8.23 The heat flow taken at the centre of the system for each simulation normalised to the local Spitzer value.

$E = \frac{1}{2}mv^2$ $\times T_{\text{cold}}$	f^1/f^0 $\times 10^{-2}$	f^2/f^0 $\times 10^{-3}$	$(\log f^0(E) - \log f^0(E+\Delta E))/\Delta E$	
			'A'	'B'
0.5	-0.31	-0.32	0.416	0.427
1.5	-1.1	-0.30	0.415	0.423
2.5	0.60	-0.30	0.414	0.421
3.5	0.87	-0.34	0.413	0.420
4.5	2.8	-0.40	0.413	0.417
5.5	4.9	-0.45	0.412	0.415
6.5	7.1	-0.49	0.411	0.414
7.5	9.3	-0.52	0.410	0.412
8.5	11.5	-0.53	0.410	0.411
9.5	13.6	-0.54	0.410	0.410
10.5	15.7	-0.54	0.410	0.409
11.5	17.7	-0.54		

Table 8.03

$E = \frac{1}{2}mv^2$	$f^1_{sp}/f^1_{\text{code}}$	f^0/f^1_{code}	f^{02}/f^1_{code}
0.5	1.2	1.0	1.0
1.5	2.1	1.0	1.0
2.5	-6.2	1.1	1.0
3.5	2.8	1.3	1.0
4.5	1.2	1.5	1.0
5.5	3.1	1.7	1.0
6.5	4.9	1.7	1.0
7.5	6.8	1.9	1.0
8.5	8.8	2.1	1.0
9.5	10.9	2.2	1.0
10.5	13.0	2.2	1.0
11.5	15.3	2.3	1.0

Table 8.04

'Maxwellian-ness' of points 'A' and 'B' on Figure 8.23. The last two columns of Table 8.03 show the deviation of f^0 from a Maxwellian for 'A' and 'B'. 'B' has clearly less of a Maxwellian structure. The heat flow for 'B' is, however, closer to the Spitzer value. The deciding factor here is the degree to which the local temperature and density gradients represent the 'effective' temperature and density gradients of the energetic heat-carrying electrons.

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8.4.3 Time-independent solution of the Vlasov-Fokker-Planck equation for a planar ablating plasma

Thermal electron energy transport in the presence of the steep temperature gradients has received much attention in recent years, and it has become apparent that the heat flow depends on its context and that a universal value for a flux limit does not exist. We report here on a self-consistent solution of the Vlasov-Fokker-Planck (VFP) equation for electrons and the hydrodynamic equations for cold ions in a planar steadily ablating plasma. The problem is made time-independent and the equation for the electron distribution is

$$v_x \frac{\partial f}{\partial x} + E \frac{\partial f}{\partial v} + u \frac{\partial f}{\partial x} - \frac{\partial u}{\partial x} v_x \frac{\partial f}{\partial v} = \frac{\partial f}{\partial t} \bigg|_c + \frac{\partial f}{\partial t} \bigg|_s$$

$$\frac{\partial f}{\partial t} \bigg|_c = \frac{1}{v^2} \frac{\partial}{\partial v} \left[v^2 \left(\frac{D_{11}}{2} \frac{\partial f}{\partial v} + Cf \right) \right] + \frac{1}{2} \frac{1}{v^2} \frac{\partial}{\partial \mu} \left[D_{\perp} (1 - \mu^2) \frac{\partial f}{\partial \mu} \right]$$

where $\mu = v_x/v$ and $\partial f/\partial t|_s$ is an externally applied heat source. The electron distribution function f is defined in a rest frame which is moving at a local velocity $u(x)$ which is set equal to the bulk plasma flow velocity. We make the following approximations:

- (1) The ions are cold and electron-ion energy exchange is omitted.
- (2) The electron distribution function can be written in the form

$f = f_0 + \mu f_1$. This is equivalent to a Legendre polynomial expansion in μ with all terms except the first two neglected. Work by Matte and Virmont (8.36) with a computer program in which higher order terms are included has shown that this is a good approximation, particularly when the heat flow is close to the regime in which the Spitzer conductivity is valid. This is confirmed by our own experience with a similar computer program (8.12), and it has also been shown to be the case in related analytic solutions of the VFP equation (8.37). Other terms in the VFP equation which make a contribution smaller in magnitude than the higher order Legendre terms are also dropped from the calculation.

(3) The collision coefficients C , $D_{\perp}$, and $D_{\parallel}$ are calculated from the isotropic part of the distribution, and are the same as those used in our previous work (8.12). They can be written as

$$D_{\parallel} = \frac{18}{Z+1} \frac{1}{\lambda_0} \frac{1}{v^3} \left[\frac{n_e kT}{m_e} - \int_v^{\infty} \frac{4\pi}{3} (u^4 - v^3 u) f_0(u) du \right]$$

$$C = \frac{9}{Z+1} \frac{1}{\lambda_0} \frac{1}{v^2} \left[n_e - \int_v^{\infty} 4\pi u^2 f_0(u) du \right]$$

$$D_{\perp} = \frac{1}{Z+1} \frac{9}{\lambda_0} \frac{1}{v} \left[(Z+1)n_e - \frac{n_e kT}{m_e v^2} - \int_v^{\infty} 4\pi \left(\frac{4uv}{3} - u^2 - \frac{u^4}{3v^2} \right) f_0(u) du \right]$$

where $\lambda_0 = 9m_e^2 / [4\pi n_e^4 (Z+1) \ln \Lambda]$

and kT/m is one third the mean square electron velocity. In the high-velocity limit, the integrals in these expressions and the term $n_e kT/m_e v^2$ tend to zero. Most of the energy transport takes place at velocities of 3 or 4 times the thermal speed v_t ($v_t^2 = kT/m_e$). We neglect these terms and made the approximation that the high velocity asymptotic forms are correct at all velocities. This overestimates the strength of the collision terms at low velocity while maintaining the correct equilibrium ratio of $D_{\parallel}/C$. Although the asymptotic approximation is good for velocities at which energy is transported, there are significant electron currents at lower velocities, and the approximation in the collision coefficients affects

the electric field required for charge neutrality. This in turn affects the heat flow. Nevertheless, the departure from the Spitzer conductivity (8.30) is determined by departure of f_0 from Maxwellian, and this only takes place at high velocities in excess of $3v_t$. Consequently, the approximate coefficients are good for our purpose provided that we compare the calculated heat flow with a 'modified Spitzer' conductivity which is also calculated from the asymptotic form of the collision coefficients. The modified Spitzer conductivity differs from the correct Spitzer conductivity by 20 per cent when $Z = 4$, and by 5 per cent when $Z = 16$, with the error decreasing as Z increases. The validity of this approximation was examined by calculating the exact coefficients and comparing them with the asymptotic approximation. The errors introduced by the approximation were found to be small for velocities at which the distribution departs from Maxwellian.

(4) It is assumed that all energy applied externally to the plasma is given to low velocity electrons ($v < v_t$) as would be appropriate for heating by inverse bremsstrahlung. In these cases considered here, heating is weak in the sense that the characteristic heating time is small compared with the electron collision time. The heating term is therefore not significant for the high velocity electrons which are active in energy transport. The input energy is given to low velocity electrons and transferred collisionally to high velocity electrons which transport it to the colder parts of the plasma. When considering energy transport it is therefore not necessary to model the energy input accurately provided that the heating is weak enough for $(\partial f / \partial t)_s$ to be negligible at high velocities. Instead, it is important to use collision coefficients which account for collisional energy transfer to those electrons which transport energy. The equations for the collision coefficients at high velocity show that this is equivalent to specifying the correct temperature which then determines the ratio $D_{\parallel}/C$. We therefore let $(\partial f / \partial t)_s = 0$, and input energy to the system by writing the coefficient $D_{\parallel}$ as $D_{\parallel} = 18 n_e \theta / [\lambda_0 (Z+1) v^3]$ and choosing θ equal to the required temperature instead of the temperature calculated by averaging over the distribution. It is found in general that the temperature calculated by averaging over the electron distribution is not exactly equal to the imposed temperature

because the omission of the term for energy input implies that the electron energy distribution is not being correctly modelled at low velocity. The difference between the two temperatures would in reality be removed by rearrangement of the electron distribution by external heating. In the case of heating by inverse bremsstrahlung this would be due to a flattening of the distribution close to zero velocity (8.38). The size of the error introduced by treating heat input in this way can be estimated from the difference between the calculated and imposed temperatures which is found at most to be a few per cent in the cases considered here. This indicates that the errors are not large and make little difference to the distribution at high velocity where the energy transport occurs. If however, heat is applied much more strongly by limiting its application to a smaller region of plasma it is found that the difference between θ and the calculated temperature is large and a full treatment including the heating term $(\partial f / \partial t)_s$ and correct collision terms is needed.

In order to solve the equation for the electron distribution we must specify the imposed temperature $\theta(x)$, the plasma flow velocity $u(x)$, and the value of f_0 at zero velocity, which is equivalent to specifying the density. The density, velocity and temperature profiles are determined by the fluid equations for mass, momentum and energy conservation. Plasma heating is included in the equation for energy conservation, and we specify the energy input arbitrarily as taking place over a distance of 10-20 electron mean free paths with a sinusoidal spatial dependence. A realistic ablating plasma cannot be truly steady-state because incoming laser light would be absorbed by inverse bremsstrahlung in a plasma ablating uniformly to infinity. However, the planar steady-state formulation with arbitrarily imposed energy input will be a good representation of a system in which the distance between the absorption and solid surfaces is small compared with other distances such as the radius of curvature of a curved critical surface. The high and low density boundary conditions can be determined in terms of shocks (8.39) which are moving away from the region of energy input, but in order to limit the size of the calculation we impose boundaries at left and right at which the derivatives are zero and hence there is no conducted heat flow. Although this is not strictly

correct, the distances of the boundaries from the active region of the plasma are large enough for the boundaries not to affect the solution noticeably. We specify the temperatures at each of the boundaries, and the Mach number and density at one of the boundaries. We choose a Mach number of 0.8 at the low density boundary. A value closer to unity would probably be more realistic but the calculation does not converge if the Mach number is close to unity. The equations are solved by an iterative process on a grid which has 64 grid-points in position and 48 in velocity. Apart from the graphics routines, the computer program is completely separate from that used in our previously published work (8.12).

The results of the calculation are given in Figure 8.24 which shows the temperature, heat flow and density profiles for a plasma with $Z = 15$ and $m_i/Zm_e = 3672$. The spatial coordinate is in units of the mean free path of an electron with energy $3kT/2$ at the hot end of the plasma. The results are compared with the Spitzer conductivity in Figure 8.24. Curves a, b and c are plots of the flux-limited Spitzer heat flow with no flux limit, and harmonically imposed flux limits of 10 per cent and 5 per cent respectively. The heat flow given by these curves is calculated from the temperature profiles $\theta(x)$ calculated by solution of the VFP equation. They show that the VFP conductivity is lower than the Spitzer conductivity near the heat source, but marginally higher on the tail of the profile. Clearly there is no one flux limit which is correct but a limit around 5 per cent appears to be appropriate near the heat flow maximum. Figure 8.25 is a plot of Q/Q_f against L/λ where $Q_f = n_e kT(kT/m_e)^{1/2}$ and $L = T/\nabla T$. The maximum value of Q/Q_f is determined by the Mach number and m_i/Zm_e , and although Q/Q_f does not exceed 3 per cent, a somewhat larger flux limit is sufficient to produce the deviation from Spitzer on account of the limit being imposed as a harmonic mean, $Q^{-1} = Q_{\text{Spitzer}}^{-1} + \alpha Q_f^{-1}$, where α is the flux limit. A second comparison of the results with the Spitzer conductivity is given in Figure 8.24 by the curve for the temperature θ_s calculated by self-consistent solution of the Spitzer conductivity with the conservation equations for mass, momentum and energy. This profile can be directly compared with the temperature calculated by solution of the VFP equation, the only difference being the way in which the heat flow

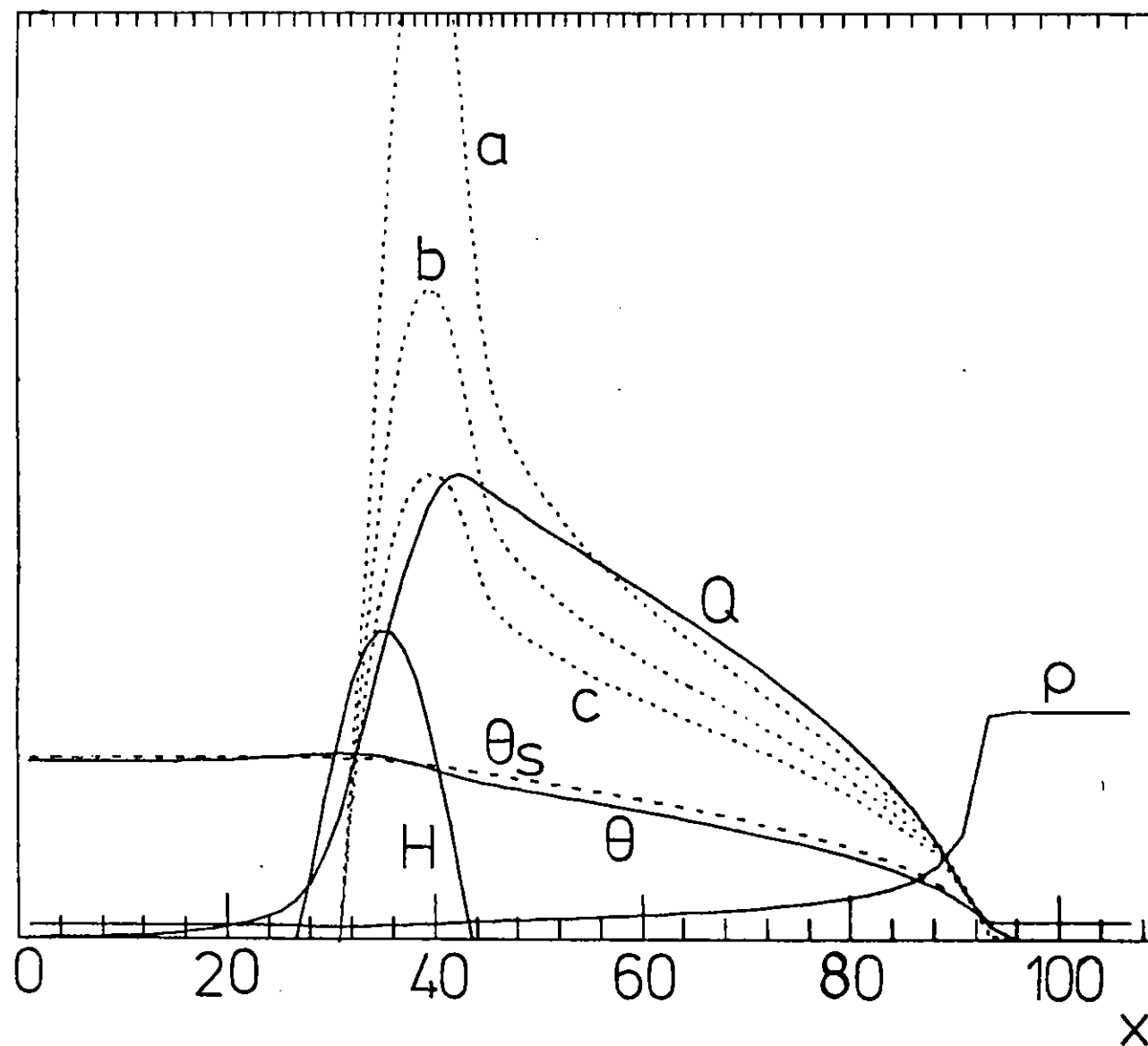


Fig. 8.24 Temperature (Θ), density (ρ), heat flow (Q), and energy input (H) as functions of position. Distance is measured in units of the mean free path of a thermal electron at the hot end of the plasma. The marks at the top of the graph indicate the positions of the grid-points. The dotted lines are the flux-limited heat flow calculated for the temperature profile Θ for (a) no flux limit, (b) a 10% flux limit, and (c) a 5% flux limit. The dashed line Θ_s is the temperature profile found by self-consistent solution of the hydrodynamic equations with the Spitzer conductivity instead of the VFP equation. The electron-ion mass ratio is $m_i/Zm_e = 3672$, and $Z = 15$.

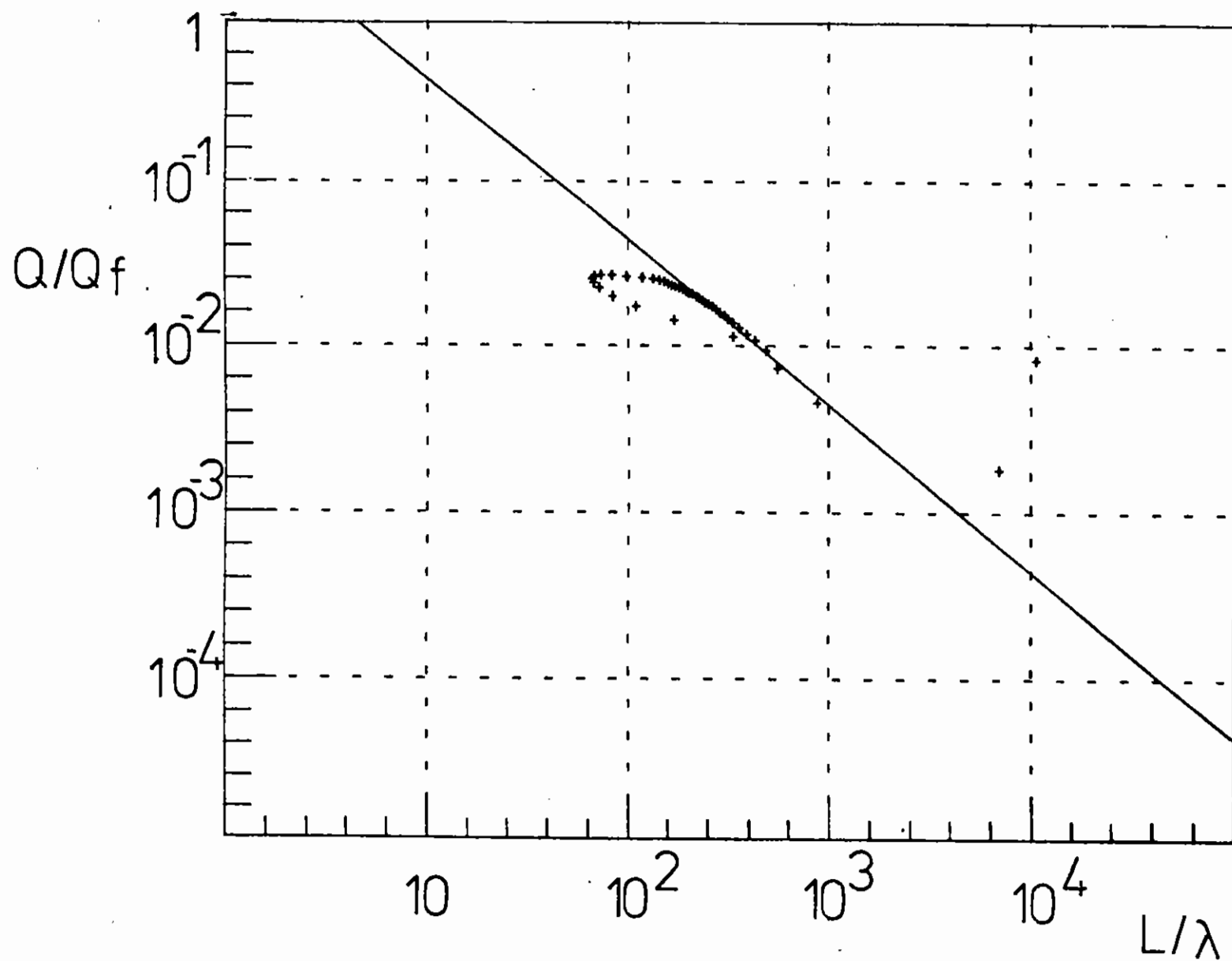


Fig. 8.25 $\log (Q/Q_f)$ plotted against $\log (L/\lambda)$ where $Q_f = n_e kT (kT/m_e)^{1/2}$, $L = T/\nabla T$ and λ is the mean free path of an electron with energy $3kT/2$. All quantities are calculated locally. Each cross represents one grid-point. The solid line is the Spitzer conductivity. The electron-ion mass ratio is $m_i/Zm_e = 3672$, and $Z = 15$.

is calculated. The difference between the VFP profile and the Spitzer profile is small. The heat flow deviates significantly from the Spitzer heat flow near the region of energy input, but the effect on the overall temperature profile is weak because the deviation only occurs over a short length of plasma, and the overall distance between the absorption and solid surfaces is little affected. The deviation from the Spitzer theory manifests itself primarily as a steepening of the temperature gradient in the region where heat is applied. The difference between the Spitzer and VFP profiles is small enough for it to be very doubtful whether this could be responsible for experimental measurements which require the use of a flux limit in a hydrodynamic computer model.

The data given in Figures 8.24 and 8.25 relate to a simulation in which the Mach number is 0.8 at the low-density boundary. The flow in a laser-produced plasma is usually divergent either because the target is spherical or because the focal spot of the laser is limited in size, in which case the Mach number could be expected to be larger. Although we cannot model higher Mach numbers with the present planar-geometry computer program, the important feature associated with a higher Mach number is the higher ratio of the plasma flow velocity to the electron thermal velocity and we can mimic this effect by changing the electron-ion mass ratio. Energy conservation implies that the maximum heat flow $Q_{\max}$ occurring on the high density side of the region of energy input is given by

$$\frac{Q_{\max}}{Q_f} = \left(\frac{Zm_e}{m_i}\right)^{\frac{1}{2}} \frac{M}{2} (5 + M^2)$$

where M is the Mach number. The ratio of the maximum heat flow $Q_{\max}$ to the free-streaming limit Q_f is increased by either increasing the Mach number or increasing the electron-ion mass ratio. Consequently, it is instructive to solve the equations with an increased electron-ion mass ratio but keeping the Mach number equal to 0.8. Figures 8.26 and 8.27 are graphs of the spatial profiles and the conductivity for mass ratio m_i/Zm_e of 632 which is equivalent to increasing the Mach number to 1.5 while keeping the mass ratio equal to 3672. The VFP conductivity differs

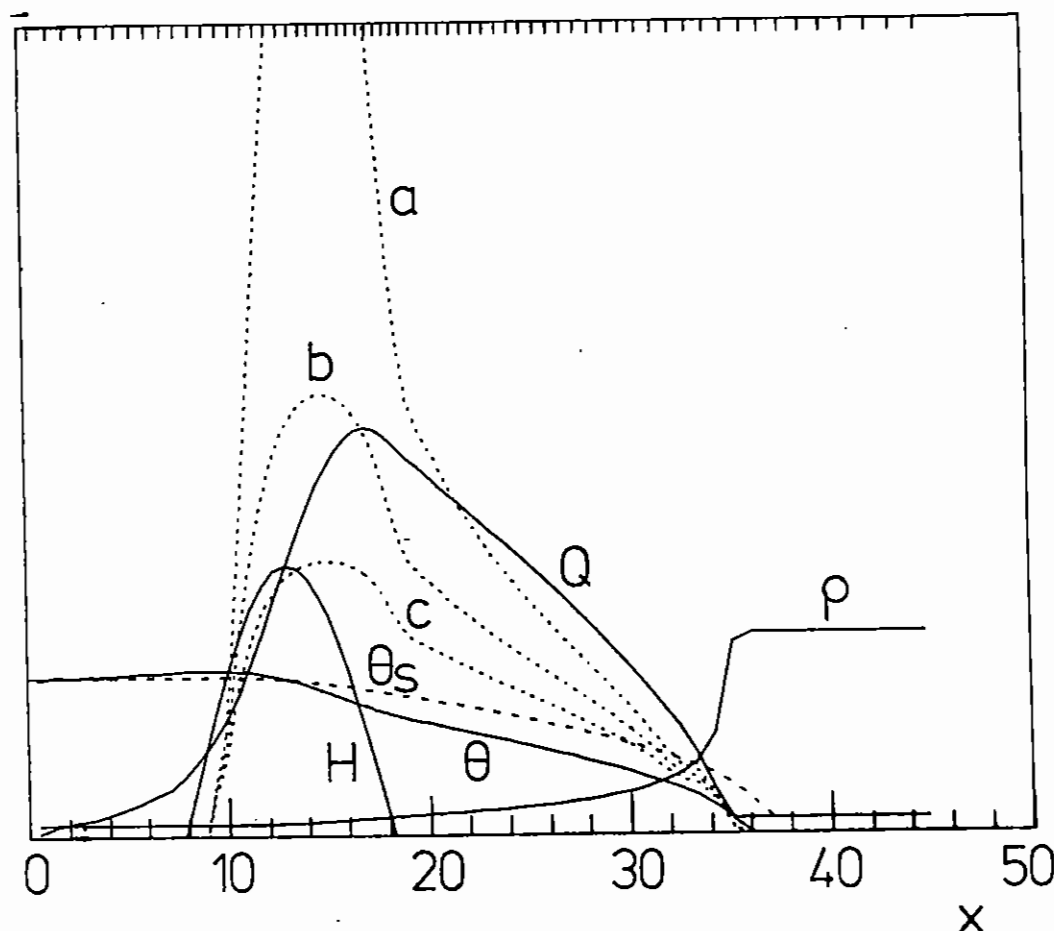


Fig. 8.26 Same as Figure 8.24, but with $m_i/Zm_e = 632$.

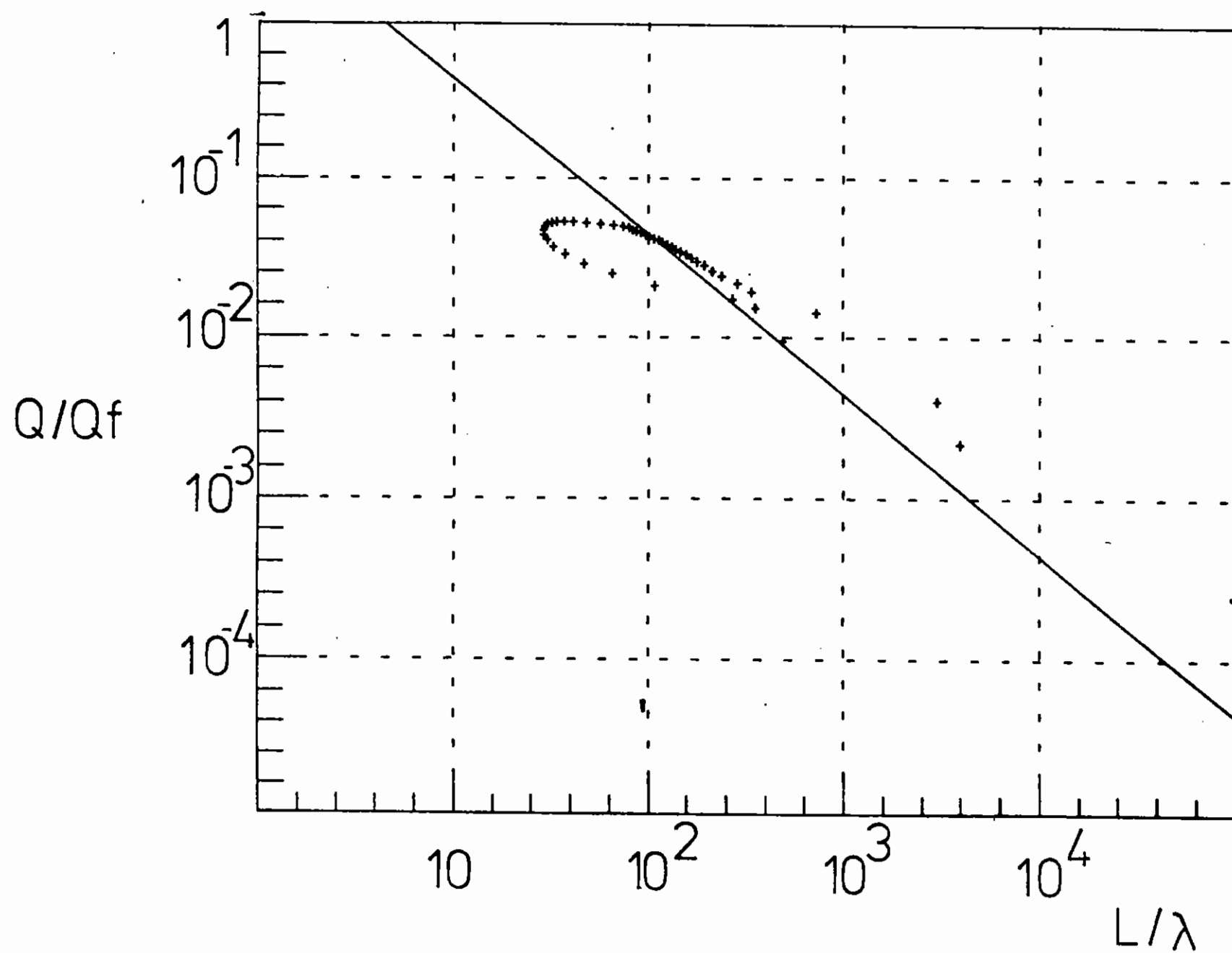


Fig. 8.27 Same as Figure 8.25, but with $m_i/2m_e = 632$.

considerably from the Spitzer conductivity. The maximum in the heat flow is commensurate with a harmonically imposed 10 per cent flux limit. To the left of the maximum in Figure 3.26, the heat flow is lower than that given by a 10 per cent limit, and to the right it is higher. When we compare the VFP temperature profile with the Spitzer profile we see once again that there is very little difference in the distance between the absorption and solid surfaces. Hence, even for an equivalent Mach number of 1.5, although a flux limit of around 10 per cent gives the best fit to the conductivity, the difference from the Spitzer conductivity occurs in such a way that measurement of the difference may be beyond the sensitivity of present experiments with laser-produced plasmas.

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8.4.4 Heat Flow Instabilities

The presence of heat flow instabilities in laser-driven plasmas may result in a wide variety of unstable phenomena (8.40). These could generally have adverse effects on the symmetry of the implosion and more specifically on the electron energy flow between critical and ablation surfaces (8.41).

Here we are concerned with magnetic field generating thermal instabilities which may grow on a time scale longer than the electron-ion collision time and whose characteristic wavelengths are larger than the electron mean-free-path. Under these conditions fluid theory is usually thought to provide an adequate description of the plasma (8.42), and therefore many instability mechanisms have been developed on the basis of such a theory (8.40). Our aim is to use a kinetic model, in an attempt to investigate the fluid predictions.

We shall start off with the assumption of no electron-electron collisions and motionless ions. The former being valid for high Z plasmas only. Since we are considering collisional electrons, a further simplification is achieved by expanding the distribution f as $f = f_0 + \underline{f}_1 \cdot \underline{v}/v$. Making use of the orthogonal properties of this expansion, the Fokker-Planck

equation becomes (8.43):

$$\frac{\partial f_0}{\partial t} + \frac{v}{3} \nabla \cdot \underline{f}_1 - \frac{1}{3} \frac{1}{v^2} \frac{\partial}{\partial v} \left(\frac{v^2 e}{m} \underline{E} \cdot \underline{f}_1 \right) = 0 \quad (1)$$

$$\frac{\partial \underline{f}_1}{\partial t} + v \nabla f_0 - \frac{e}{m} \underline{E} \frac{\partial f_0}{\partial v} - \underline{\Omega} \wedge \underline{f}_1 + v \underline{f}_1 = 0 \quad (2)$$

where the collision frequency is $\nu = n_1 Y / v^3$, $Y = 4\pi (Ze^2/m)^2 \ln \Lambda$, $n_1 (=n/Z)$ being the ion number density, $\ln \Lambda$ the Coulomb logarithm, $\underline{E}$ the electric field, and $\underline{\Omega}$ the electron cyclotron frequency.

In order to couple the fields in the Fokker-Planck equation we need Faraday's and Ampere's Law,

$$\frac{\partial}{\partial t} \underline{\Omega} = - \frac{e}{m} \nabla \wedge \underline{E} \quad (3)$$

$$\nabla \wedge \underline{\Omega} = \frac{4\pi e j}{mc} = - \frac{(4\pi e)^2}{3mc} \int_0^\infty dv v^3 \underline{f}_1 \quad (4)$$

where, by neglecting the displacement current we have restricted ourselves to describing phenomena on a scale larger than a Debye length.

Linearizing equations (1) to (4) by assuming an equilibrium and transverse perturbation of the form: $\underline{E} = (0, 0, E_z^0(x)) + (\underline{E}_x^1(x), 0, \underline{E}_z^0(x)) \exp i(kx - \omega t)$, $\underline{\Omega} = (0, \Omega_y^1(x), 0) \exp i(kx - \omega t)$, $f = f_0^0(z) + f_z^0(z) v_z / v + (f_x^1(x) + f_x^1(x) v_x / v + f_z^1(x) v_z / v) \exp i(kx - \omega t)$, we obtain,

$$-i\omega f_0 + \frac{ikv}{3} f_x^1 = 0 \quad (5)$$

$$v f_x^1 + ikv f_0 - \Omega_y^1 f_z^0 - \frac{e}{m} \frac{E_x^1}{v} \frac{\partial f_0}{\partial v} = 0 \quad (6)$$

$$v f_z^1 + v \frac{\partial}{\partial z} f_0 - \frac{e}{m} \frac{E_z^1}{v} \frac{\partial f_0}{\partial v} - \frac{e}{m} \frac{E_z^0}{v} \frac{\partial f_0}{\partial v} = 0 \quad (7)$$

$$i\omega\Omega_y^1 = -\frac{e}{m} \left(ikE_z^1 - \frac{\partial}{\partial z} E_x^1 \right) \quad (8)$$

$$ik\Omega_y^1 = -\frac{(4\pi e)^2}{3mc} \int_0^\infty dv v^3 f_z^1 \quad (9)$$

where the perturbation wavelength $\lambda(=2\pi/k)$ is assumed to be smaller than a characteristic scale length L , in accordance with the local approximation. Following the collisionality condition, imposed on the electrons, we have neglected from equations (6) and (7) the inertia term $i\omega f_z^1$ in comparison with $v f_z^1$, and shall require that $\lambda \gg \lambda_{th}$, where $\lambda_{th} = (2T/m)^{1/2} n_1^{-1}$ is the mean-free-path of an electron of energy T . Lastly, we have neglected the electric field contribution from equation (5), in view of the fact that j and Ω are first order quantities.

Since we have omitted the displacement current in our analysis, we need to evaluate the longitudinal electric field E_x^1 by means of the quasi-neutrality condition $\nabla \cdot j = 0$, i.e. $j_x^1 = 0$. Thus, substituting (5) into (6) and taking the v^3 moment of the resultant expression for f_x^1 ,

$$f_x^1 = \left(\frac{-3i\omega}{k^2 v^2 - 3i\omega v} \right) \left(\Omega_y^1 f_z^0 + \frac{eE_x^1}{m} \frac{\partial f_z^0}{\partial v} \right) \quad (10)$$

we obtain $eE_x^1/m = \phi \Omega_y^1$,

where $\phi = \left[-\int_0^\infty dv v^3 f_z^0 / (k^2 v^2 - 3i\omega v) \right] \left[\int_0^\infty dv v^3 \frac{\partial f_z^0}{\partial v} / (k^2 v^2 - 3i\omega v) \right]^{-1}$.

Inserting E_x^1 back into equations (5)-(8) we arrive, after lengthy algebraic manipulations, at an integro-eigenvalue equation in $f_z^1: f_z^1 + \chi(\omega, k, f^0) \int_0^\infty dv v^3 f_z^1 = 0$. This has standard solutions of the form $f_z^1 = \text{constant} \times \chi$ with eigenvalues ω given by $1 + \int_0^\infty dv v^3 \chi = 0$. Expressing the dispersion relation in terms of dimensionless parameters, we have,

$$1 + \frac{1}{3k^2 L_s^2} \int_0^\infty dv v^6 \left[\left(\frac{v \partial}{\partial z} - \epsilon^0 \frac{\partial}{\partial v} \right) \left(\frac{v^5}{v^5 - n} \right) \left(\frac{F_z^0}{v} + \frac{\phi}{v} \frac{\partial F_z^0}{\partial v} \right) + iW \frac{\partial F_z^0}{\partial v} - \frac{\partial \phi}{\partial z} \frac{\partial F_z^0}{\partial v} \right] = 0 \quad (11)$$

where,

$$W = \omega/v_{th}, \quad \phi = \phi/v_{th},$$

$$V = v/v_{th}, \quad L_s = c/\omega_p \lambda_{th},$$

$$K = k\lambda_{th}, \quad d^n/dz^n F = (4\pi/n)v_{th}^3 d^n/dz^n f,$$

$$\epsilon^0 = eE_z^0/mv_{th} v_{th}, \quad n = 3i\omega v_{th}/k^2 v_{th}^2,$$

with $v_{th} = (2T/m)^{1/2}$ being the thermal velocity and $v_{th} = n_1 Y/v_{th}^3$ the collision frequency of an electron of energy T , and $\omega_p = (4\pi n e^2/m)^{1/2}$ the plasma frequency.

In the limit of small n and by assuming a 'Spitzer' type equilibrium distribution, equation (11) reduces to,

$$W_I = A_1 \frac{1}{L_t^2} + A_2 \frac{1}{L_n L_t} + A_3 \frac{\Delta}{L_t^2} - \frac{\sqrt{\pi}}{8} k^2 L_s^2 \quad (12)$$

where, $L_{n,t} = \lambda_{th} d/dz \ln(n,t))^{-1}$, $\Delta = \lambda_{th} L_t^2 d^2/dz^2 T$, and $W = W_r + iW_i$. Coefficients A_1, A_2, A_3 are listed on table 8.0.

In the fluid approach one would have initially taken appropriate velocity moments of the Fokker-Planck equation, evaluated the transport coefficients, and then linearized the resultant set of equations. Tidman and Shanny (8.44), and then later Hirao and Ogasawara (8.45), have developed the thermomagnetic instability theory based on Braginskii's fluid equations and obtained a dispersion relation equivalent to (12). For better comparison between kinetic and fluid results we have also developed a fluid type instability theory by taking velocity moments of our linearized transport equations (5)-(7). On table 8.05 we list

	A_1	A_2	A_3
Braginskii Coefficients	9.099	-0.3896	6.066
Our Coefficients	9.814	1.908	6.543
Kinetic Theory	14.87	-1.453	6.543

Table 8.05

	α_0	β_0	γ_0	$\frac{\beta_0^1}{\sigma_0}$	$\frac{\gamma_0^1}{\sigma_0}$
Braginskii Coefficients	0.2949	1.521	12.471	9.126	106.5
Our Coefficients	0.2945	1.500	13.581	9.844	172.7

Table 8.06

coefficients A_1 , A_2 , A_3 , based on both fluid models.

By comparing the fluid coefficients on table 8.05 we observe different values for A_2 . The reason for this arises from a discrepancy between our transport coefficients and those derived by Braginskii (8.42). We have listed them on table 8.06, where we notice in particular that the Right-Leduc heat flux coefficient obtained by us is a factor of 1.6 higher than Braginskii's.

A comparison between fluid and kinetic results, on the other hand, indicates a much larger discrepancy. This comes about because of the assumption of no electron-electron collisions in our kinetic model. Thus, rather than being always close to thermal equilibrium, as assumed in fluid theory, the electrons are able to diffuse in space much more freely, as governed by the diffusion equation $\frac{\partial f_0^1}{\partial t} + \frac{v}{3} \nabla \cdot f_0^1 = 0$. Clearly the isotropic part of the perturbed distribution f_0^1 will not necessarily be of Maxwellian form. A very similar effect has been identified by Bell (8.37) in the context of ion waves. Strictly, the validity of our kinetic results should be verified by comparing the magnitude of the e-e collision term with other terms in the Fokker-Planck equation.

We have solved the full dispersion relation (11) numerically and compared it with the equivalent fluid results. In Figures 8.28 and 8.29 we plot W_1 against $\Lambda (= \lambda/\lambda_{th})$ for $V_n \cdot \nabla T < 0$ and $V_n \cdot \nabla T > 0$ respectively. As observed, the unstable modes obtained from kinetic theory are more pronounced for $V_n \cdot \nabla T < 0$ and for decreasing electron heat flow q . We also notice that the discrepancy between kinetic and fluid results diminishes as $\Lambda > 10^3$.

In conclusion, we have rederived the fluid transport coefficients in the limit of high Z and small magnetic fields, and found the Right-Leduc heat flux coefficient to be 60% larger than obtained by Braginskii. We have also found that for perturbation wavelengths shorter than about 1000 electron mean-free-paths, the non-thermal nature of the electrons tends to drive the thermomagnetic instability harder than is predicted by fluid theory. This effect being most significant in the regions where $V_n \cdot \nabla T < 0$

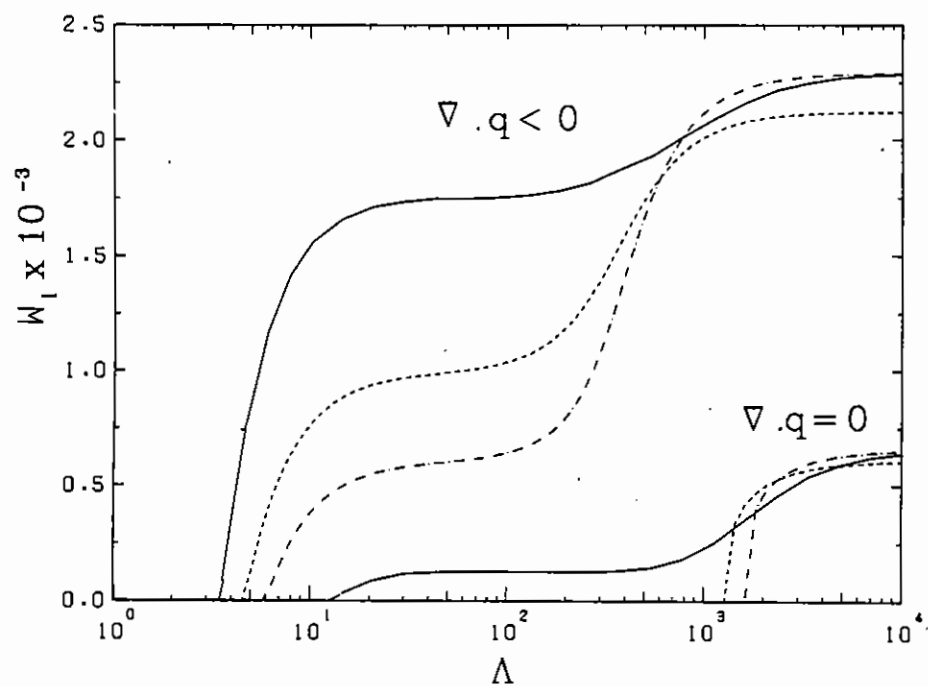


Fig. 8.28 Plots of W_I versus Λ for $L_n = 50$, $L_t = -100$, $L_s = 0.05$ and $\Delta = -5/2, 0$ (ie $\nabla \cdot q = 0$ and $\nabla \cdot q < 0$ respectively). Kinetic results are represented by the full line, fluid results with our coefficients by the dash-dotted line, and fluid results Braginskii's coefficients by the dashed line.

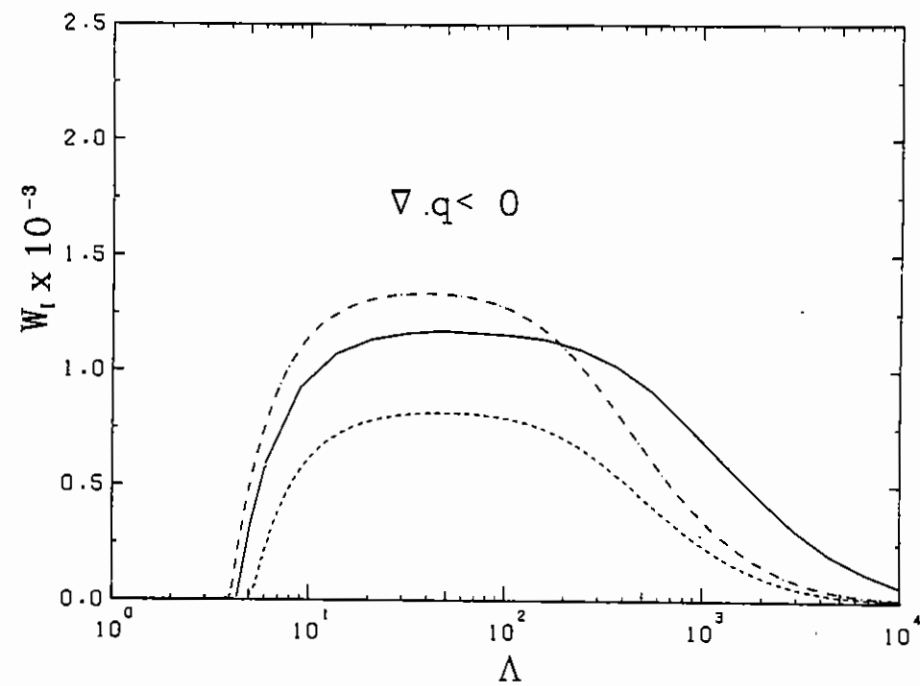


Fig. 8.29 Same as Figure 8.28 but with $L_t = 100$.

and $\nabla \cdot \underline{q} < 0$, as would be typical in the overdense plasma.

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8.4.5 Radiation Transport with a Planckian Distribution

Further investigation has been made of the radiation energy equation (8.46) when the radiation field is Planckian. The relevant source terms in the equation are due to electron-radiation energy exchange, which we will compare with electron-ion energy exchange, and radiation thermal conductivity, which we investigate when the radiation temperature varies from that of the background plasma. Both of these terms may be derived from the equation of transfer, which in the usual notation is

$$\frac{1}{c} \frac{\partial I_{\nu}}{\partial t} + \underline{\Omega} \cdot \nabla I_{\nu} = \kappa_{\nu}^1 (I_{\nu p} - I_{\nu}) \quad (1)$$

This is appropriate for material in LTE, and where scattering can be ignored. We proceed in the usual manner, taking the zeroth and first angular moments of (1). Combining these in the diffusion approximation yields

$$\frac{\partial U_{\nu}}{\partial t} + \nabla \cdot \left(\frac{-c}{3\kappa_{\nu}^1} \nabla U_{\nu} \right) = c \kappa_{\nu}^1 (U_{\nu p} - U_{\nu}) \quad (2)$$

where U_{ν} is the spectral energy density. $U_{\nu p}$ is its equilibrium value, tied as it were, to the temperature of the plasma. After integrating over frequency, we may identify the second term on the LHS of (2) as the radiation thermal conduction term, whilst on the RHS we have the electron-radiation energy exchange term. The former may be shown to be

$$\frac{-c}{3} \nabla \cdot \left[\lambda_R \nabla \left(\frac{4\sigma T_e^4}{c} \right) \right] \quad (3)$$

where we have defined the Rosseland mfp in terms of the non-equilibrium Planckian U_{ν} , and hence T_e , by

$$\lambda_R \equiv \kappa_R^{-1} = \frac{\int_0^{\infty} (\kappa_{\nu}^1)^{-1} \frac{\partial U_{\nu}}{\partial T_e} d\nu}{\int_0^{\infty} \frac{\partial U_{\nu}}{\partial T_e} d\nu} \equiv \frac{\int_0^{\infty} (\kappa_{\nu}^1)^{-1} W_R(\nu) d\nu}{\int_0^{\infty} W_R(\nu) d\nu} \quad (4)$$

It is customary at this stage to define λ_R in terms of the equilibrium density $U_{\nu p}$, and hence T_e . To determine the significance of this we have evaluated both possibilities for comparison. Since κ_{ν}^1 must be specified we have chosen the simplest example, that of a fully ionised gas, where $\kappa_{\nu} \propto \nu^{-3}$. Thus we may represent κ_{ν}^1 (corrected for stimulated emission) in terms of the reduced frequency $u = h\nu/kT_e$ as $A(1 - \exp(-u))/u^3$ where $A = A(T_e, N_1, N_e)$. For our radiation Planckian however, our weighting function W_R is given by

$$\frac{15}{4\pi^4} \alpha^5 u^4 e^{\alpha u} / (e^{\alpha u} - 1)^2 \quad (\alpha = T_e/T_r)$$

We now see that the customary choice uses $\alpha = 1$. We shall therefore compare $\lambda_R(\alpha)/\lambda_R(\alpha = 1)$ for a fully ionised gas. This is plotted in Figure 8.30. For $0.1 \alpha < 10$, the ratio varies over six orders of magnitude, ie adoption of the equilibrium value seems inadmissible with large temperature variations.

The electron-radiation energy exchange term will utilise the definition of Planckian mean opacity, κ_p , viz

$$\kappa_p = \frac{\int_0^{\infty} \kappa_{\nu}^1 U_{\nu p} d\nu}{\int_0^{\infty} \kappa_{\nu}^1 U_{\nu p} d\nu}$$

or

$$\int_0^{\infty} \kappa_{\nu}^1 U_{\nu p} d\nu = \frac{4\sigma T_e^4}{c} \kappa_p$$

This is defined in terms of the equilibrium density. It is tempting to write the frequency integrated RHS of (2) as $4\sigma\kappa_p(T_e^4 - T_r^4)$, since then the electron-radiation exchange term would be proportional (inter alia) to $(T_e - T_r)$, analogous to the electron-ion case. The fallacy lies in the definition of κ_p which cannot simultaneously be defined in terms of $U_{\nu p}$ and U . We shall now show that the temperature ratio $\alpha = T_e/T_r$ and not

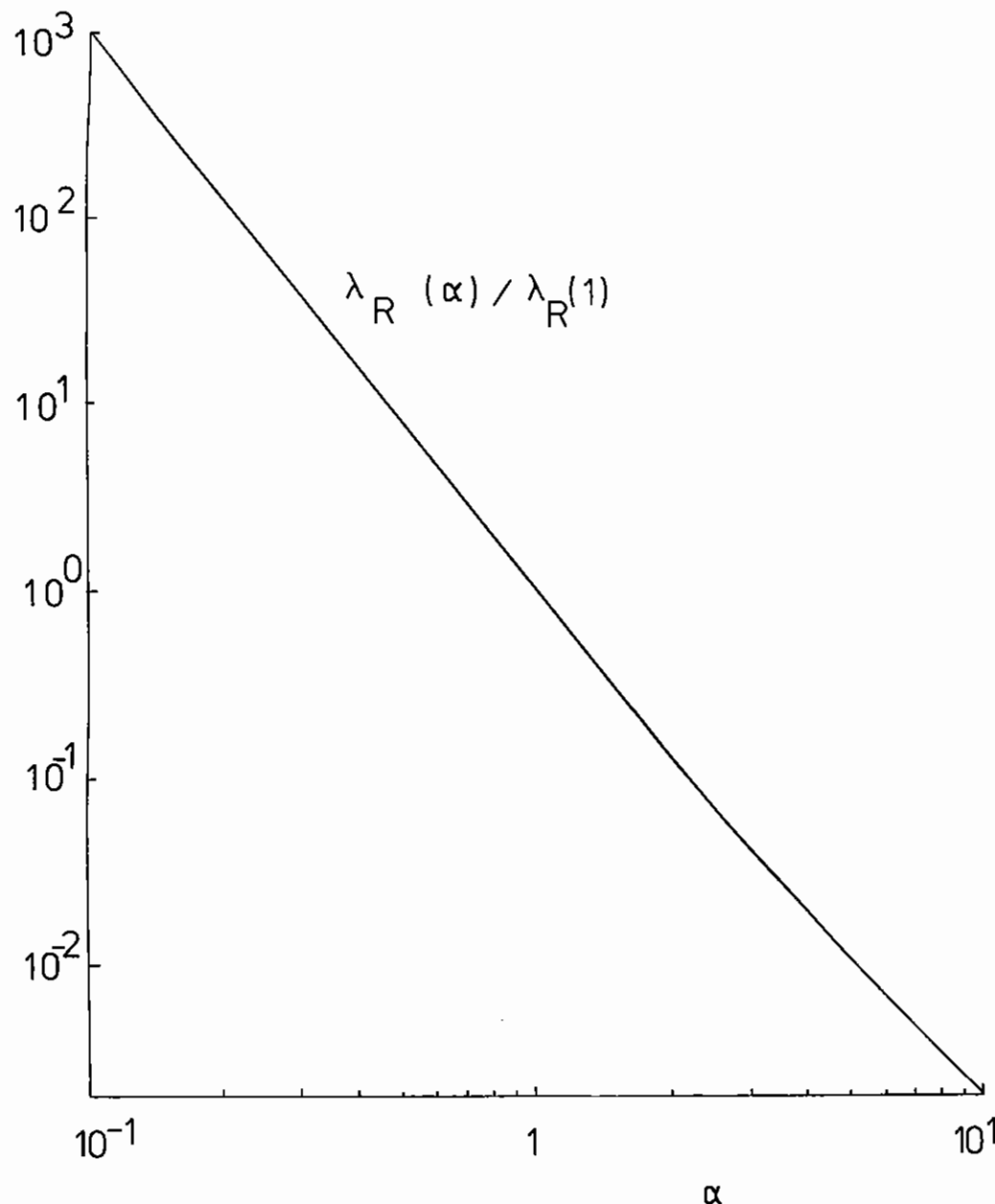


Fig. 8.30 The ratio of $\lambda_R(\alpha)/\lambda_R(1)$ (see text).

temperature difference ($T_e - T_r$) is the relevant parameter in the exchange term (again for a fully ionised gas). In terms of the reduced frequency, the frequency integrated RHS of (2) becomes

$$B \int_0^{\infty} \frac{(1 - e^{-u})}{u^3} \left[\frac{u^3}{e^u - 1} - \frac{u^3}{e^{\alpha u} - 1} \right] du$$

where $B = A 8\pi h/c^2$, with A as above. In passing we note that the stimulated emission factor is crucial to ensure the convergence of the integral. The integral may be replaced by $-(\psi(1/\alpha) + \gamma)/\alpha$ where ψ is the Psi/Digamma function. γ is Euler's constant (see Figure 8.31). Curiously, as α increases, ie $T_e \gg T_r$, the graph indicates that the electron to radiation exchange approaches a limiting rate. Conversely there seems to be no limit to the rate at which energy can be lost to the electrons from the radiation. As expected there is no exchange when $\alpha = 1$, ie $T_e = T_r$, but the temperature difference is nowhere involved.

In conclusion, however, it must be conceded that the above results, both for the Rosseland mfp variation with T_r and the electron-radiation exchange rate, will not be as dramatic as they seem since the LTE requirement will restrict their application to cases where T_r is near T_e .

A C White (Glasgow University)

8.4.6 LTE Opacity Calculations

The radiative opacity of hot materials is important for the calculation of radiative preheat in laser targets. A computer code, originally written by T R Carson (St Andrews) to calculate the LTE radiative opacity of stellar material, is being adapted to produce multigroup opacities for use in laser target calculations.

The opacity calculation uses the average atom model and for any given temperature and density it firstly calculates the Thomas-Fermi spherically symmetric potential. This is used to obtain both bound and free one-electron wave functions from a numerical solution of the Schroedinger

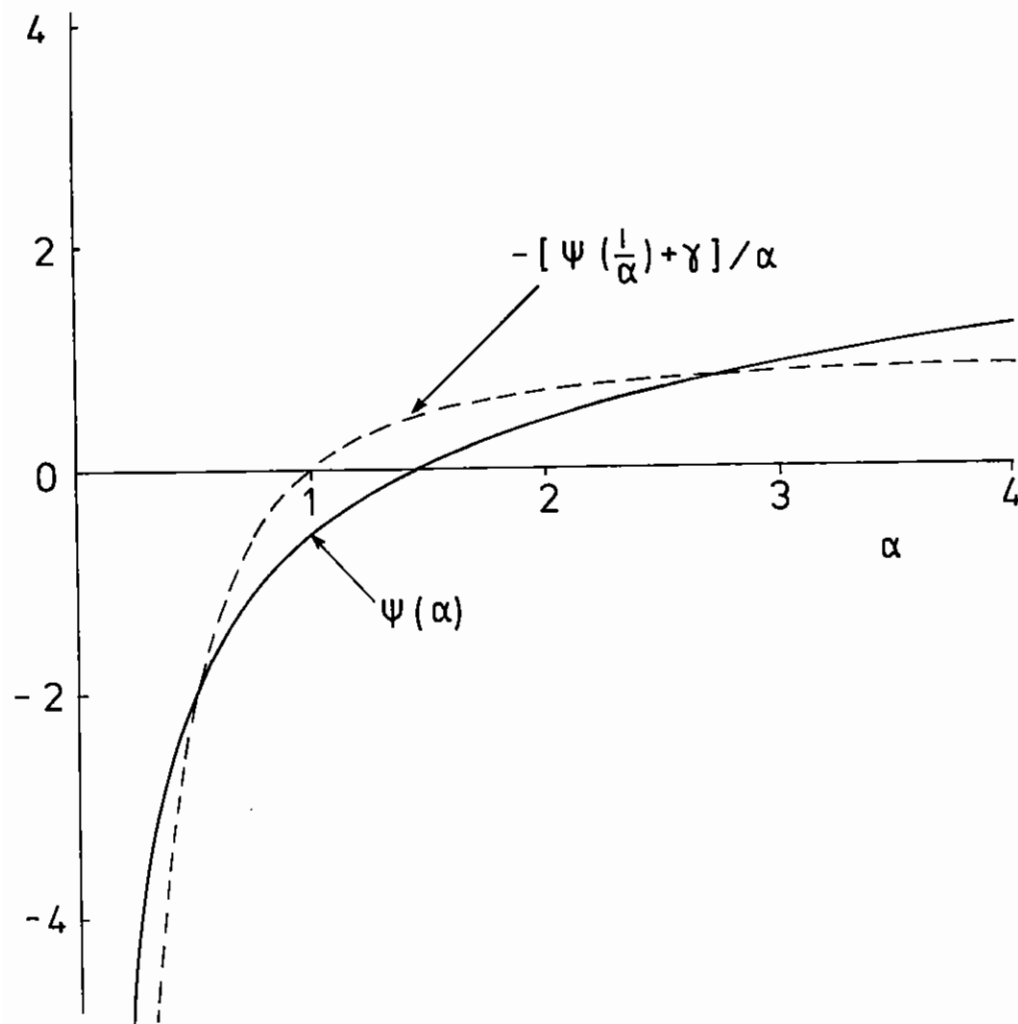


Fig. 8.31 (a) The psi/digamma function.
 (b) $-(\psi(1/\alpha) + \gamma)/\alpha$. $\alpha = T_e/T_r$.

equation. The probability of occupancy of each one-electron state comes from Fermi-Dirac statistics. The code includes bound-bound, bound-free and free-free contributions to the opacity which are calculated from matrix elements of the one-electron wave functions together with the average electron occupancies.

It is hoped that the code will produce multigroup data for the microballoon shell materials glass and plastic as well as the fill gasses Freon, DT and Neon. Initially the code is being used to calculate single-group Rosseland opacities to allow comparison with existing stellar opacity data.

T R Carson (St Andrews), D J Nicholas, S J Rose (RAL) and A C White (Glasgow)

8.4.7 The Calculation of Radiation Preheat in Laser Driven Implosions

There are two areas for which radiation transport needs to be considered in laser produced plasmas. The first is the calculation of the X-ray spectrum including details of line shapes and line intensity ratios for comparison with spectroscopic observations. The second area, which is the subject of this work, is the calculation of energy flow by X-ray photons which can affect the behaviour of a laser target both by radiation cooling of the hot core and radiative preheat. To assess the effect of radiative preheating, the transfer of energy by radiation has been included in the 1D Lagrangian hydrodynamics code MEDUSA.

In a laser-heated plasma, the region of maximum X-ray emission occurs just ahead of the critical surface and energy is radiated both outwards and inwards to the cooler absorbing region. The diffusion approximation is therefore questionable because of its assumption of isotropy of the radiation field. The model adopted (which with some modifications is that described in the previous annual report) differentiates between photons travelling into the target and those travelling outwards. In each case it is assumed that radiation is isotropic in the half space. The photon mean free path depends strongly on the photon energy, in general the higher the

photon energy, the longer the mean free path. To account for this in an approximate manner, the calculations split into 30 groups and each group is transported separately; inelastic scattering is expected to be negligible and is ignored. Only photon absorption and emission is considered. Attenuation of the flux through each cell is taken to be exponential. Because the cool, dense region for which absorption is thought to be most important is expected to be in LTE, the absorption opacity is calculated from an LTE average-atom model due to Pitzer. However the X-ray emitting region is not expected to be in equilibrium and the emission from each cell is calculated assuming a coronal steady state reduced by a correction factor determined by comparison with experimental measurements of X-ray spectra and with X-ray conversion efficiency measurements. This correction factor is thought to arise because the ions do not spend sufficiently long in the heating region for the ionic populations to reach a steady state.

The effect of radiative preheat on the implosion of thick-walled plastic microballoons can be seen from Figures 8.32, 8.33 and 8.34. Figure 8.32 shows that although the motion of the inside surface of the plastic shell is relatively unaffected by radiative preheat, the outside of the shell is more strongly affected, the 0.1 g/cc contour appears at a larger radius when radiation preheat is included. This is because the density profile has been modified by the inclusion of radiative preheat (Figure 8.33). This may have considerable implications for the interpretation of backlighting implosion experiments. The effect of radiation transport on the electron temperature through the shell can be seen in Figure 8.34 where radiation heats the inside surface of the target.

S J Rose and R G Evans (RAL)

8.4.8 Some Numerical Aspects of NLTE Radiation Transfer Calculations

There exists a large body of work relating to Non-LTE radiation transfer in stellar atmospheres (8.47)(8.48). Previous work connected with the CLF has used the equivalent two level atom (ETLA) approach (8.49) to perform time-independent calculations (8.50). Additionally time-integrated

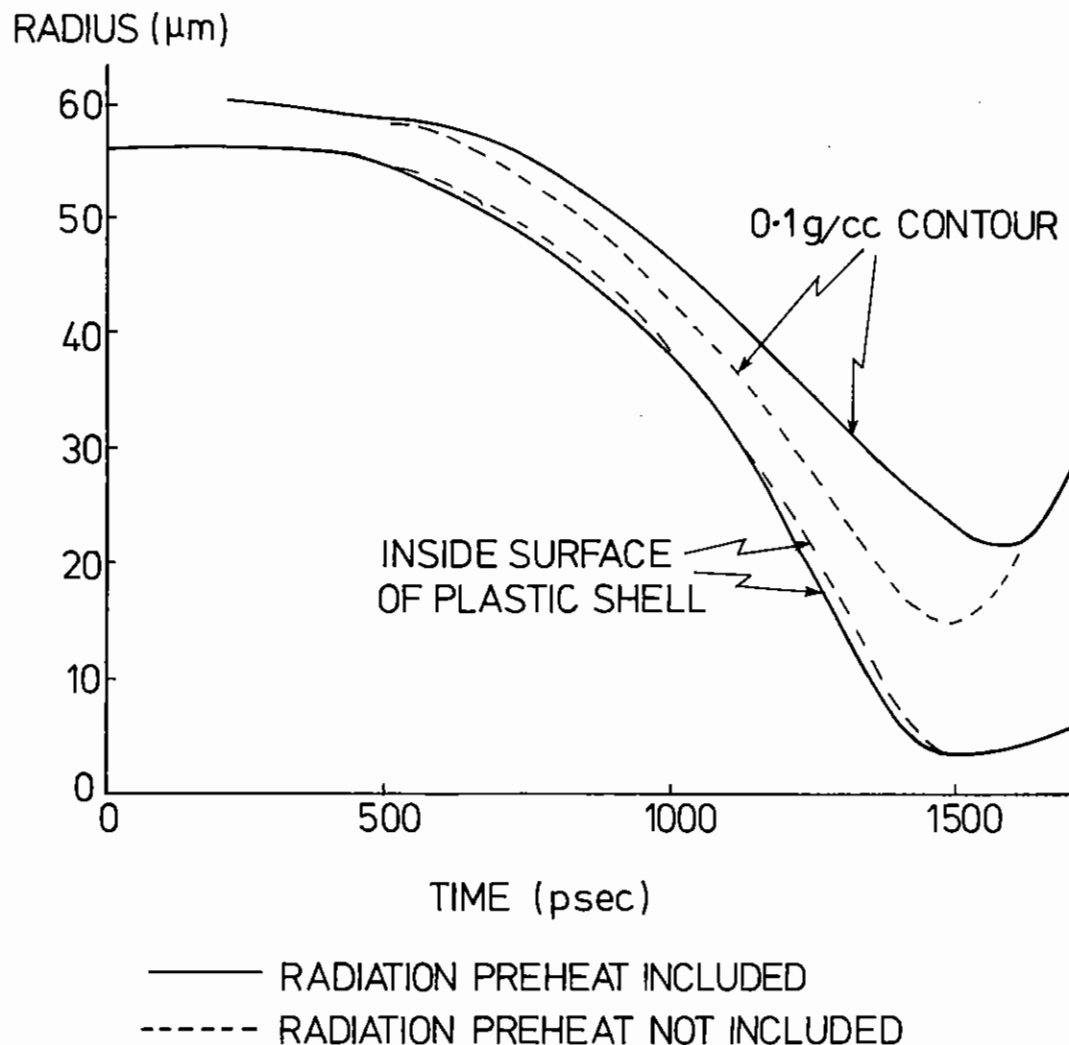


Fig. 8.32 The effect of radiation transport on the implosion of a microballoon (123 μm diameter; 5.8 μm thickness; $8 \times 10^{13} \text{ W cm}^{-2}$ incident green light; 1.2 ns pulse length).

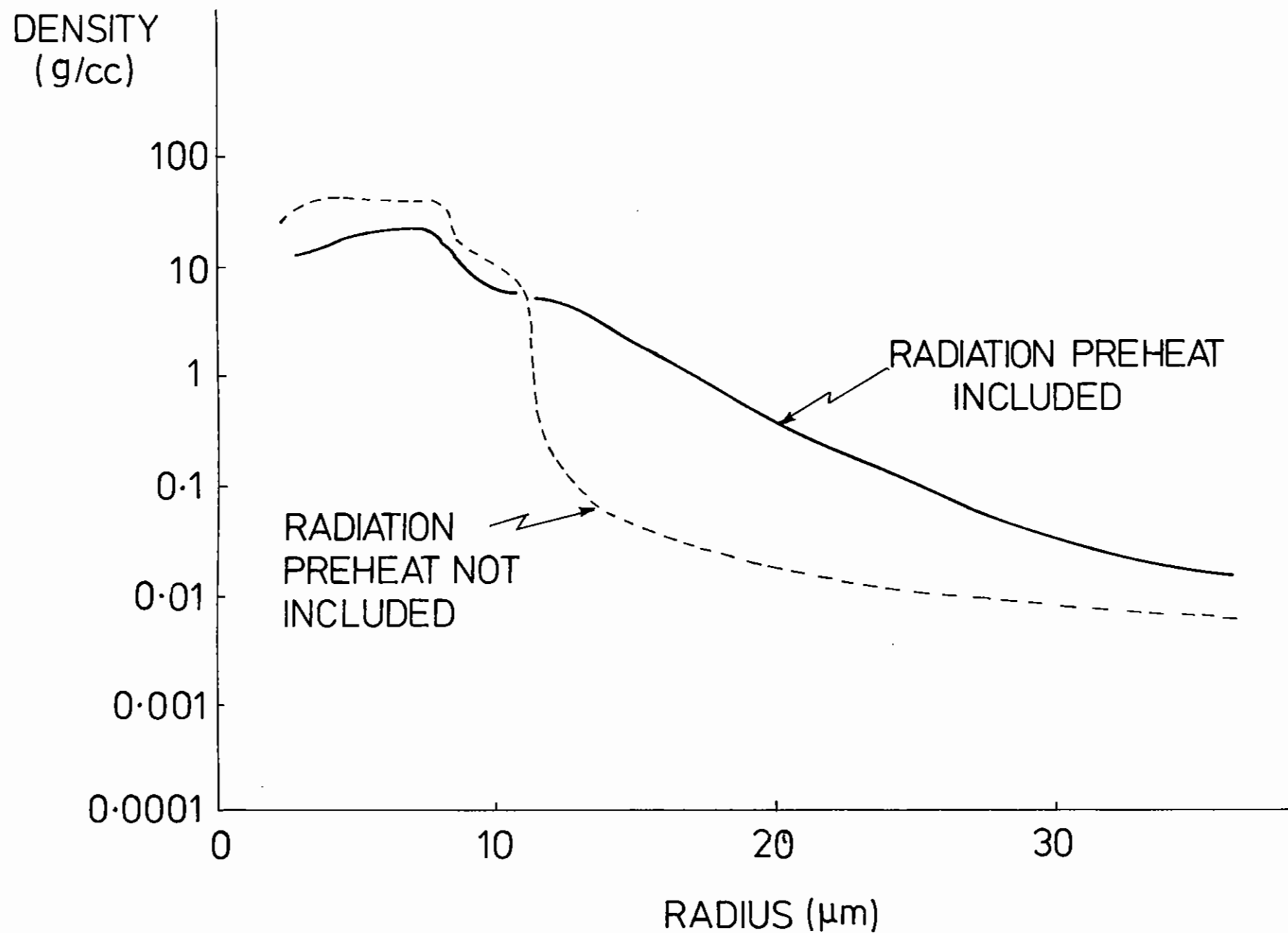


Fig. 8.33 Density of the plastic shell at 1.5 ns.

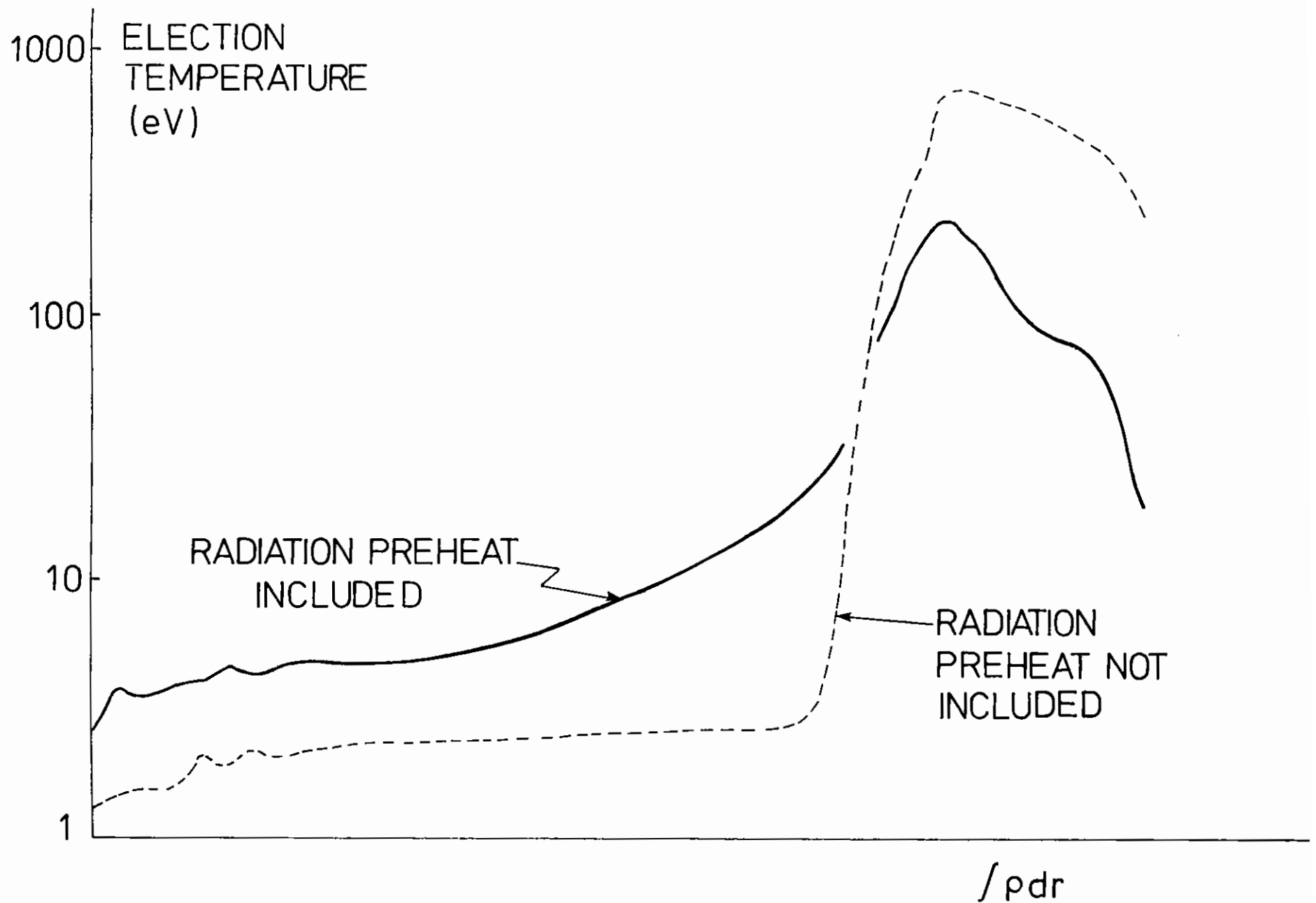


Fig. 8.34 Electron temperature of the plastic shell at 1.0 ns.

spectra have been synthesised by summing the results of time-independent calculations (8.51).

Three aspects of performing radiation transfer calculations in laser heated targets have been considered.

- (a) The inclusion of time-dependent radiation transfer and rate equations within the ETLA approach.
- (b) What is the most efficient way of solving the radiative transfer equation in the fairly collisional plasmas of moderate optical depth encountered in laser heated targets?
- (c) Simulating radiation transfer in media where population inversion exists.

Time dependent radiation transfer

The time dependent form of the radiation transfer equation is:

$$\frac{1}{c} \frac{\partial I}{\partial t} + \mu \frac{\partial I}{\partial x} = -\chi(t) I + \eta(t)$$

Time dependent effects arise because of the $\frac{1}{c} \frac{\partial}{\partial t}$ term and the time dependence of the opacity and emissivity (χ and η). If fully implicit differencing is used in both the transfer and atomic rate equations then it is straightforward to obtain a finite difference scheme using the ETLA approach. The computational work required is comparable to that needed to solve the time independent problem. Details are given elsewhere (8.52).

Efficiency of Calculations

Two numerical approaches to the solution of time independent radiative transfer - kinetic equilibrium calculations exist. The Radiation Physics Group at the US Naval Research Laboratories have employed Λ iteration. This approach involves iterating between a formal (χ and η given) solution

of the transfer equation and a solution of the atomic rate equations. It may be adapted to model time dependent problems in a straightforward manner (8.53).

Other investigators (8.50) (8.54) have used the ETLA approach. A provisional optical depth scale is constructed for each transition. Assuming that the rates from levels other than the pair being considered can be taken as fixed it is then possible to derive an expression for the source function (η/χ) which is linear in the radiation field. Simultaneous solution of the radiation transfer equation for all frequencies within the spectral line of interest allows some account of the effect of the radiation field on the source function to be made implicitly. The atomic rate equations are then solved and the whole process is iterated.

The ETLA approach converges in fewer iterations than Λ iteration but the computational cost per iteration is far greater for the ETLA method.

The ratio of the costs of Λ iterations to ETLA iteration is highly problem dependent. If (collisional de-excitation rate/spontaneous radiative rate) > 0.1 or optically moderately thin (probability of escape > 0.1), Λ iteration has a very considerable advantage.

The advantage which Λ iteration has may be increased and extended to less collisional and more optical thick plasmas if the iterative process can be accelerated. The Chebyshev semi-iterative scheme (8.55) has been successfully applied to radiative transfer problems (8.56). Additionally the method developed by Reed (8.57) has been used to find regimes where the synthetic method (8.58) should be useful (8.59).

Radiative Transfer in Lasing Media

Alley et al (8.54) have studied gain in externally pumped X-ray laser schemes. They encountered numerical difficulties in treating the radiative transfer in inverted transitions. These arise for two reasons. Firstly, because the Feautrier (u, v) form of the radiation transfer

equation is inherently unstable if the opacity is negative. Secondly, because the idea of using a provisional optical depth scale is no longer useful. These problems may be avoided by (a) using the first order form of the radiative transfer equation and (b) accounting for the change in opacity and emissivity of the lasing transitions to first order in the change in the radiation field.

It is then possible to solve the radiative transfer equation in a manner which implicitly accounts (to first order) for the change of χ and η due to the radiation field. This approach has been successfully tested on some simple problems.

D J Bond (Imperial College)

8.4.9 1982 CECAM Workshop on 'The Flux Limiter and Heat Flux Instabilities'

This three-week workshop, attended by over 20 specialists from Europe and the USA, covered three areas; one-dimensional non-linear heat flow, heat flow instabilities, and now of increasing importance, two dimensional heat flow. This last subject is relevant to thermal smoothing and lateral heat flow, but involves the spontaneous generation of high magnetic fields ($\sim$ several megagauss) which considerably modifies electron motion and transport. Here a two-dimensional fluid code (based on linear transport theory) was compared with a collisionless particle-in-cell code which both showed rapid lateral propagation of fast electrons with energy deposition far away from the laser focal spot region. The role of the advective Righi-Leduc heat flow was quantitatively assessed.

A new code, ANTHEM, developed by R Mason, was implemented at the Workshop: since this includes collisions in a self-consistent electro-magnetic particle code it will be able in principle to describe non-linear two-dimensional heat flow and heat flow instabilities and could be important for future understanding of the subject. In the area of heat flow instabilities a new approach involving the em perturbation of the Fokker-Planck equation in which the equilibrium describes a one-dimensional non-

linear heat flow was developed and is the subject of continuing work and collaboration. This model should put the subject of heat flow instabilities on a firm footing.

M G Haines (Imperial College)

8.4.10 The CECAM Workshop on 'Non-LTE Radiative Transfer'

The first workshop on non-LTE (non Local Thermodynamic Equilibrium) radiative transfer was held at CECAM between 20 September and 1 October 1982. The intention was to bring together those interested in the subject, to identify problems and to decide on future area of research.

Following a series of informal presentations of participants' work, discussion groups decided that benchmark problems in non-LTE radiative transfer were desirable and two such benchmarks were subsequently defined. The first consisted of a hydrogen plasma with low ion density ($2.5 \times 10^{16} \text{ cm}^{-3}$) and low temperature (0.5 eV) typical of a low temperature laboratory plasma source. The second benchmark, more typical of laser-compressed plasmas was that of a high ion number density (10^{21} cm^{-3}) and high temperature (1 keV) Argon plasma. In each case details of energy levels and transition rates to be used in the benchmark calculations were agreed. It is hoped that several participants will tackle these problems, that comparisons will be made at a future meeting and that the result of this collaboration will be agreed benchmarks.

At present non-LTE radiative transfer calculations require detailed data on energy levels and rates concerning all of the principal species in the plasma. At present only low-Z systems have been considered in this way. At conditions of interest in typical plasmas high Z systems cannot be treated in this manner, both because of the vast amount of data required and because most of that data is not known. The use of the average atom model to overcome this problem was discussed and this posed intriguing questions concerning the role of the average atom model in non-LTE radiative transfer. It was felt that this topic should be discussed at a future meeting.

Even for low-Z materials treated in current non-LTE radiative transfer calculations it was felt that some relevant data was not available and it was agreed that participants would endeavour to enlist the aid of theoretical atomic physicists in an attempt to rectify this.

R W Lee (Imperial College) and S J Rose (RAL)

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APPENDIX

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