

RAL-84-049

Science and Engineering Research Council

Rutherford Appleton Laboratory

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

**Annual Report to the Laser
Facility Committee 1984**

RAL-84-049

Laser Division

Rutherford Appleton Laboratory

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SCIENCE AND ENGINEERING RESEARCH COUNCIL

CENTRAL LASER FACILITY

ANNUAL REPORT TO THE LASER FACILITY COMMITTEE 1984

PREFACE

The report covers the work done at or in association with the Central Laser Facility (CLF) during the year ended 31 March 1984. The three main parts of the programme of the CLF are covered in three self contained sections, namely Section A, "The Glass Laser Scientific Programme", Section B, "The Ultra Violet Radiation Facility Scientific Programme", and Section C, "High Power KrF Laser Development". Section A presents the work of the six specialist scientific groups studying applications of the high power Nd glass laser. Section B reports the multidisciplinary applications of the excimer and excimer pumped frequency tunable lasers of the Ultra Violet Radiation Facility. The research work reported in both sections A and B is by University groups, often with collaboration from CLF staff and quite commonly with international collaboration, and is the result of scheduled operations of the CLF's lasers. Section C describes R&D work mainly by CLF staff on a high power KrF laser which may provide a new user facility in the future.

There were some notable achievements in the scientific programme. The Nd glass laser in its new 1kJ configuration was used to drive implosions with six beams of second harmonic green light. Optimised conditions for high density compression were chosen on the basis of much earlier study of factors affecting implosions. Sophisticated diagnostics, including the first measurement of the energy loss of thermonuclear protons, were used to diagnose the compressed density and density x radius. The results showed that compressed densities up to 10 g cm^{-3} were achieved with polymer shell targets and density x radius of $10^{-2} \text{ g cm}^{-2}$ in glass shell targets, in good agreement with computer modelling assuming perfect sphericity. These experiments are at the forefront of international study of compression by

direct laser irradiation (Section A2.2). Another notable experiment with collaboration from Los Alamos National Laboratory used the infra red output of the six beam laser to make a detailed study of heat flow in spherical geometry, casting new light on an internationally topical question about the steepness of the thermal conduction front by a novel X-ray streak spectroscopic method (Section A1.5).

The use of laser produced plasmas as uniquely intense sources of both X-rays and thermonuclear particles has been an area of great achievement this year with exciting results in time resolved X-ray diffraction (Section A4.6) and EXAFS (Section A4.4). Application to the study of technologically important processes in surface layers of semiconductors is particularly promising.

A new aspect of work with the glass laser has been the strengthening of links with the Max Planck Institut fur Quantenoptik in West Germany through a collaboration agreement with current joint work on several problems including the development of a target release system for unsupported spherical targets.

The European Conference on Laser Interaction with Matter (XVI ECLIM) was organised in 1983 by Imperial College and was an important scientific event with some 26 papers from groups associated with the CLF (Section D)

The scope of the work of the Ultra Violet Radiation Facility has continued to expand. The major achievement of the year is perhaps the development of UV phase conjugate reflection (Section B3.2) and UV photoablative etching techniques (Section B4) to a point where an independently funded project has been approved by the SERC Engineering Board and the Department of Trade and Industry to study the use of UV lasers for technologically important developments in photolithography.

Further diversification of the uses of the UVRF is evidenced by new studies of biological interest (Section B2) while 'mainstream' photochemical applications have been vigorously pursued with particularly good results in resonance Raman spectroscopy and photochemistry of small molecules (Section B1).

Finally it is gratifying to report that the high power KrF laser Sprite is currently the most powerful operational UV laser in the World.

A significant non scientific event was the retirement in June 1983 of Alan Gibson, who was an exceptionally respected and popular Division Head of Laser Division, which resulted in my call to a role where the previous act is difficult to follow ...

M H Key

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INTRODUCTION

M H Key and P T Rumsby

The Central Laser Facility (CLF) is funded by the Science Board of the Science and Engineering Research Council and in the year covered by this report, April 1983 to April 1984, its allocation was £k2200.

The Science Board has provided the facility for research whose major themes are:

- A High Density Plasma Physics
- B High Intensity Laser Plasma Interactions
- C Energy Transport in Laser Produced Plasmas
- D Laser Driven Implosions
- E X-ray Lasers
- F Other Applications of Laser Produced Plasmas
- G High Power UV Gas Laser Development
- H Applications of Short Pulse Frequency Tunable Lasers
- I Ultra Violet Lithography

The programme is managed by the Laser Facility Committee (LFC) whose membership is given in Table 1. The LFC delegates responsibility to the Rutherford Appleton Laboratory for the operation of the Facility.

The Laser Division of RAL was allocated 50 man years of effort which was deployed in the organisational structure shown in Table 2.

Staff costs and overheads used 48% of the allocation while facility running and minor development costs used 42%. 10% of the total allocation was spent on the Phase I programme of capital facility development (discussed in more detail in Sections A6 and C below). The major facilities provided by the CLF and a brief indication of their principal scientific applications is given in Table 3.

The Nd glass laser is the major facility and about 62% of resources was used to support its programme of research by UK Universities described in Section A.

Development of high power UV gas lasers is an in house project which is aimed at providing future facilities for University users. About 14% of resources are committed to this work (Section C).

Multidisciplinary scientific application of excimer lasers and excimer pumped dye lasers by UK Universities is catered for by the relatively new and small Ultra Violet Radiation Facility (using 11% of resources) which had its first full year of operations in the reporting year (Section B).

A target preparation service (8%) and a small group giving theoretical support by maintaining computer codes (5%) completes this description of the CLF.

Use of the CLF by UK Universities is free of charge to those successful in a peer review process. Research grants of typically three years duration covering expenditure in University Laboratories and research assistants' salaries may be obtained from the LFC by application to SERC. LFC Grant holders are entitled to bid for time on the CLF's lasers through a procedure described below.

A few per cent of the time may also be allocated to bids for short trial experiments from University researchers working in the Science Board area who do not hold LFC grants.

The SERC encourages use of the CLF by industry or other organisations on the basis of payment for the full cost of laser time. Such use is limited to not more than 10% of available laser time.

International collaboration in the research programme is particularly welcomed, either through collaborative participation in approved experiments of UK University groups, or through agreements negotiated with the CLF and SERC.

Selection and scheduling of University experiments using the Nd glass laser is via proposals to one of six informal specialist scientific groups which anyone may join. Each has a University Chairman and an RAL secretary (Table 4). Each group places its proposals in order of priority prior to the preparation of a draft schedule (by the programme coordinator and group secretaries). The draft schedule, which includes as many of the high priority proposals as possible (typically < 50% of proposals), is then discussed and modified at a meeting of the more formal Glass Laser Scientific Programme and Scheduling Committee (GLSPSC). All University staff members involved in the programme have voting rights, together with the group secretaries and programme coordinator. The Committee is chaired by the Division Head of Laser Division and is formally advisory to him.

A similar procedure applies to use of the UVERF but the UVERF scheduling meeting has no scientific sub groups and proposals are brought to it directly.

Scheduled experiments are supported by the operating budget of the CLF which covers travel and subsistence costs for University users and minor engineering work and other costs of experiments.

The experimental programme is linked to theoretical work in Universities which is supported by grants from the LFC and coordinated by the Theory and Computation group (Section A5).

There is a large and diverse community of University researchers participating in the programme. They number about 130 staff and research students. The majority are physicists but there has been a recent rapid growth of participation from chemists (30) and biologists.

MEMBERSHIP OF THE LASER FACILITY COMMITTEE 1983/84

Professor D D Burgess (Chairman)
Imperial College of Science
& Technology
London

Dr C L S Lewis
Department of Pure & Applied
Physics
The Queen's University of Belfast

Dr J E Allen
Department of Engineering Science
University of Oxford

Dr A Maitland
Physics Department
St Andrew's University

Dr E H Evans
Department of Biology
Preston Polytechnic

Dr H T Miles
Department of Physics
University College of Swansea

Dr D C Hanna
Department of Physics
Southampton University

Professor J P Simons
Department of Chemistry
University of Nottingham

Dr J D Kilkenny
Imperial College of Science
& Technology
London

Dr I J Spalding
Culham Laboratory
UKAEA

Professor E W Laing
Department of Natural Philosophy
University of Glasgow

Table 1

LASER DIVISION, RAL

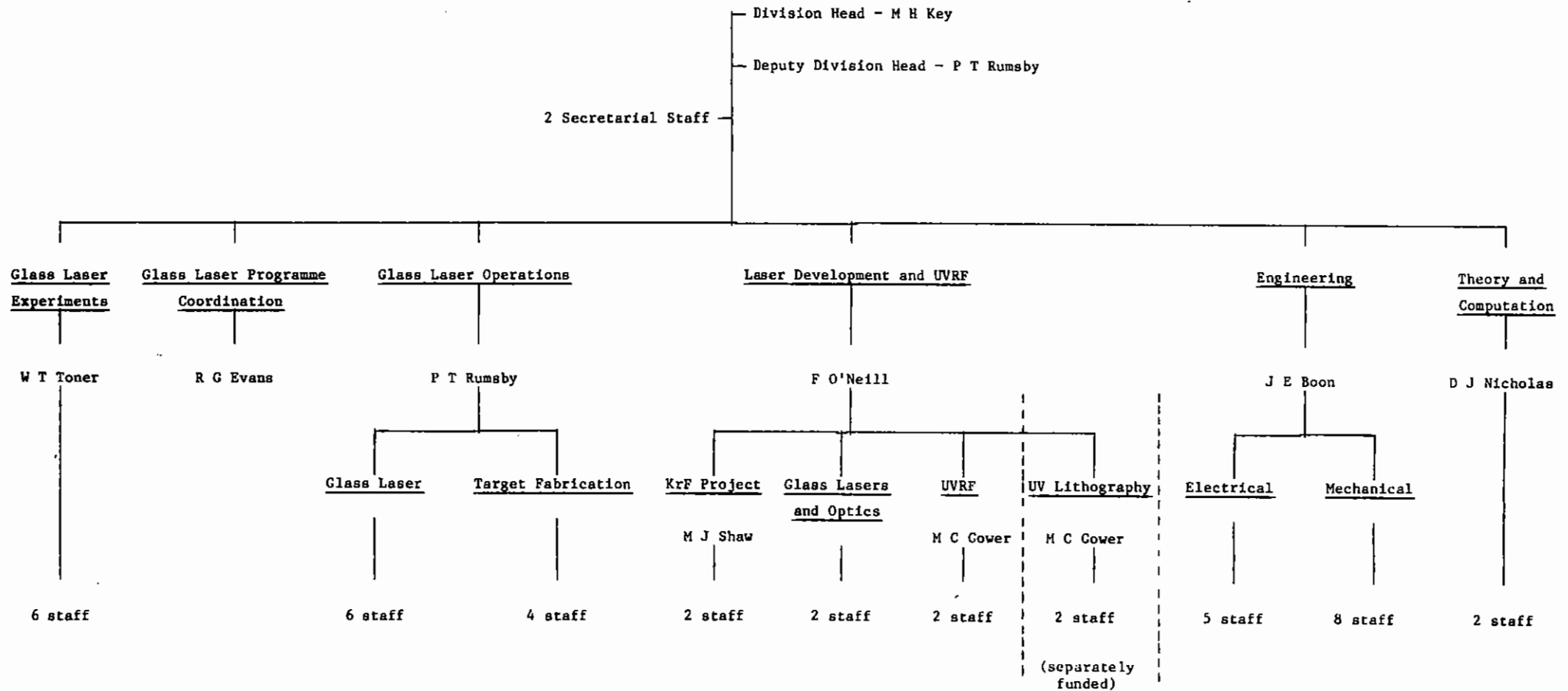


Table 2

Table 3 Major Facilities at the CLF

Nd glass laser facility	Ultra - violet radiation facility	High power KrF laser Development
VULCAN	UVRF	SPRITE
Plasma Physics	Photo chemistry	High power UV laser research
Applications of intense X-ray and particle sources	Photo biology Materials processing UV lithography Non linear optics Plasma diagnostics (Future VUV facility)	(Future facility utilizing high X-ray brightness UV laser produced plasma)
X-ray laser research		
Target Preparation Service		
Theoretical support service		

Table 4

Groups of the GLSPSC	Chairman	Secretary
Laser plasma Interaction and Energy transport	Mr T P Hughes	Mr W T Toner
Laser plasma Compression and Dynamics	Dr T A Hall	Dr A J Cole
Dense plasma and radiation physics	Prof D D Burgess	Dr S J Rose
XUV lasers and application of laser produced plasmas	Prof G J Pert	Dr R W Eason
Theory and Computation	Dr R A Cairns	Dr A R Bell
Glass Laser Facility Operations and development	Dr M H Key	Mr J E Boon
(Scientific programmer coordinator	R C Evans)	

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SECTION A GLASS LASER SCIENTIFIC PROGRAMME

INTRODUCTION

R G Evans (RAL)

In the period covered by the annual report the VULCAN glass laser has operated reliably and without any major changes in the laser and target areas. With the imminent arrival of the twelve beam illumination geometry in the new Target Area West there has been a surge of activity as users have sought to establish the baseline performance of their experiments in the existing six beam geometry. This has led to a successful series of implosion experiments using six $0.53 \mu\text{m}$ beams to drive stable, low aspect ratio ablative implosions to densities in excess of 10 g cm^{-3} and core sizes less than a tenth of the initial target radius. Transport and ablation in spherical geometry was the subject of an international collaboration with Los Alamos National Laboratory and is an area in which important theoretical and experimental advances have been made at the CLF. The measurement of charged particle energy loss rates in hot dense plasmas has been very successfully extended to a wider range of target parameters than previously and is now being used as a diagnostic of target parameters in other experiments.

Rayleigh Taylor instabilities on the outside of spherical targets have been observed for the first time to complement the earlier data on planar targets and the Essex X-ray framing camera is now being used for this work.

An important advance in high density plasma physics has been the diagnosis of shock velocity, particle velocity and density change in a colliding planar shock, together with the change in the X-ray absorption spectrum near the K-edge during shock transit.

The study of laser plasma interactions has continued its trend to using preformed, controllable target plasmas using the Imperial College Z pinch and exploding foils. The Thomson scattering data on heat flow has been extended to show structure on the ion feature in the presence of large heat

flow, while exploding foils have given important new data on the role of plasma wave propagation in producing the $3\omega_0/2$ spectrum. An important part of this has been a collaboration on the plasma theory with the University of Helsinki.

"Other" applications of laser produce plasmas have continued to grow with EXAFS being joined with high resolution X-ray imaging for biological purposes and a study of laser annealing by time resolved Bragg X-ray diffraction from the sample surface. Slow but steady progress in X-ray laser research is being made with an improved understanding of C fibre hydrodynamics.

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Editor W T Toner

A1 LASER PLASMA INTERACTION AND ENERGY TRANSPORT

A1.1 Introduction

T P Hughes (Essex) and W T Toner (RAL)

In the past year the group has continued to attack long standing problems in its field with increasingly sophisticated techniques. Substantial progress was made in our understanding of $3/2 \omega_0$ emission (Section A1.2) and thermal transport (Section A1.6) and clear evidence for SRS was observed in long scale length plasmas at moderate irradiances (Section A1.3). Visitors from Helsinki, Los Alamos and Ottawa made substantial contributions to the work.

Temporally and spectrally resolved measurements of forward and backscattered $3/2 \omega_0$ from thin foil targets are discussed in A1.2. At early times, the spectra are very broad and double peaked and pulsate on a timescale of about 30 ps. On the disappearance of the critical density surface for ω_0 the pulsations cease and the widths of the spectra decrease markedly. Although the backscattered blue and forward scattered red peaks are both reduced, they do not entirely disappear. The data are analysed in terms of a model of the two plasmon decay instability which can account for the continued presence of both peaks in forward and backscattered $3/2 \omega_0$ through refraction of the plasmons in the inhomogeneous plasma. The abrupt cessation of pulsations when the plasma goes underdense is taken as strong evidence in support of models which ascribe pulsations in $2 \omega_0$, $3/2 \omega_0$, ω_0 and $\omega_0/2$ emission to the effects of unstable fluid flow across the ponderomotive pressure jump at critical density.

In section A1.3 it is demonstrated that the thin foil target technique can be used to produce the long density scale lengths required for stimulated Raman scattering to dominate over the two plasmon decay instability at moderate irradiance. Very broad backscattered spectra are observed, centred at about 9000 Å with widths between 1500 Å and 3000 Å. An unsuccessful attempt to observe Thomson scattering from thin foil plasmas is described in A1.4.

The time resolved X-ray spectroscopy of layered targets in spherical

geometry discussed in A1.5 gives very clear evidence of a sharply defined heat front advancing at a rate which is consistent with a nearly classical thermal transport model. Measurements of absorption, magnetic fields, ion blow-off, expansion profile and X-ray continuum emission yield a comprehensive picture of the plasma conditions and place important constraints on the theoretical modelling. The disagreement between this work and results from LLE Rochester is discussed.

Section A1.6 is a progress report on the study of high gradient transport in a laser heated Z-pinch plasma. Time resolved spectra of both electron and ion features were obtained containing clear evidence of ion acoustic turbulence and also exhibiting spectral asymmetries.

A1.2 $3 \omega_0/2$ Harmonic Emission from Laser Irradiated Thin Foils

V Aboites, T P Hughes, E McGoldrick, S M L Sim (Essex), S Karttunen (Technical Research Centre of Finland) and R G Evans (RAL)

A1.2.1 Abstract

Temporally and spectrally resolved $3 \omega_0/2$ harmonic emission have been obtained from 1.05 μm laser irradiated thin foil targets which become underdense during the laser pulse of 600 - 900 ps. The results show a change in the spectrum of both the backscattered and forward scattered harmonic light coincident with the disappearance of the n_c surface.

The results are analysed by considering the phase matching conditions for the generation of the $3 \omega_0/2$ light, including, for the first time, the effects of plasma wave refraction in the inhomogeneous plasma.

A1.2.2 Introduction

The observation ((A1.1) to (A1.4)) of harmonic emission at the frequency $3 \omega_0/2$ implies the generation of $\omega_0/2$ plasma waves near the quarter critical density surface, $n_c/4$. These $\omega_0/2$ plasmons can be generated by either the two plasmon decay instability (A1.5) or by stimulated Raman scattering (A1.6). In an inhomogeneous plasma the two plasmon decay instability is

expected to be the dominant mechanism due to its lower threshold. The spectrum of the $3\omega_0/2$ harmonic emission arising from the coupling of the two plasmon decay waves with incident or reflected electromagnetic waves has been calculated by several authors ((A1.3) and (A1.7)) to be double peaked with wings shifted to both the blue and the red of the nominal $3\omega_0/2$ position.

The (ω, k) matching conditions (A1.7) for the production of $3\omega_0/2$ radiation from the decay waves in the two plasmon instability are such that the blue wing of the spectrum can only be generated propagating away from the incident laser beam. The blue wing can be observed in backscatter only if it subsequently reflects from its own critical surface at $9n_c/4$.

The experiment described here was designed to investigate the contributions to the $3\omega_0/2$ spectrum from reflected electromagnetic waves. As well as the reflection of the $3\omega_0/2$ light there is ω_0 light reflected from the n_c surface which can combine with $\omega_0/2$ plasma waves at $n_c/4$. If a thin foil is irradiated with a moderately long laser pulse the foil will explode and its peak density will reduce with time. When the peak density falls below $9n_c/4$ the reflected $3\omega_0/2$ light will disappear and when the peak density falls below n_c the effects of reflected ω_0 light will disappear. If the $3\omega_0/2$ light is produced solely by the two plasmon decay process then all $3\omega_0/2$ emission should cease when the peak density falls below $n_c/4$. In this experiment the $3\omega_0/2$ light emitted in forward and backscatter is spectrally and temporally resolved and transmitted ω_0 light gives a time marker as the peak density in the foil target drops below n_c . Changes in the $3\omega_0/2$ spectrum are expected with the blue wing in backscatter disappearing at the same time as the $9n_c/4$ surface and the (less certain) effects of reflected pump radiation disappearing at the time of ω_0 "burnthrough".

A1.2.3 Experimental Arrangement

The experimental layout is shown in Figure A1.1. The incident light beam ($\lambda_0 = 1.053 \mu\text{m}$, $\tau_{\text{FWHM}} = 600$ to 900 ps), containing energies on target of up to 100 J, was focused by an $f/1.5$ aspheric doublet lens onto thin foils of various materials and thicknesses. The backscattered and forward scattered

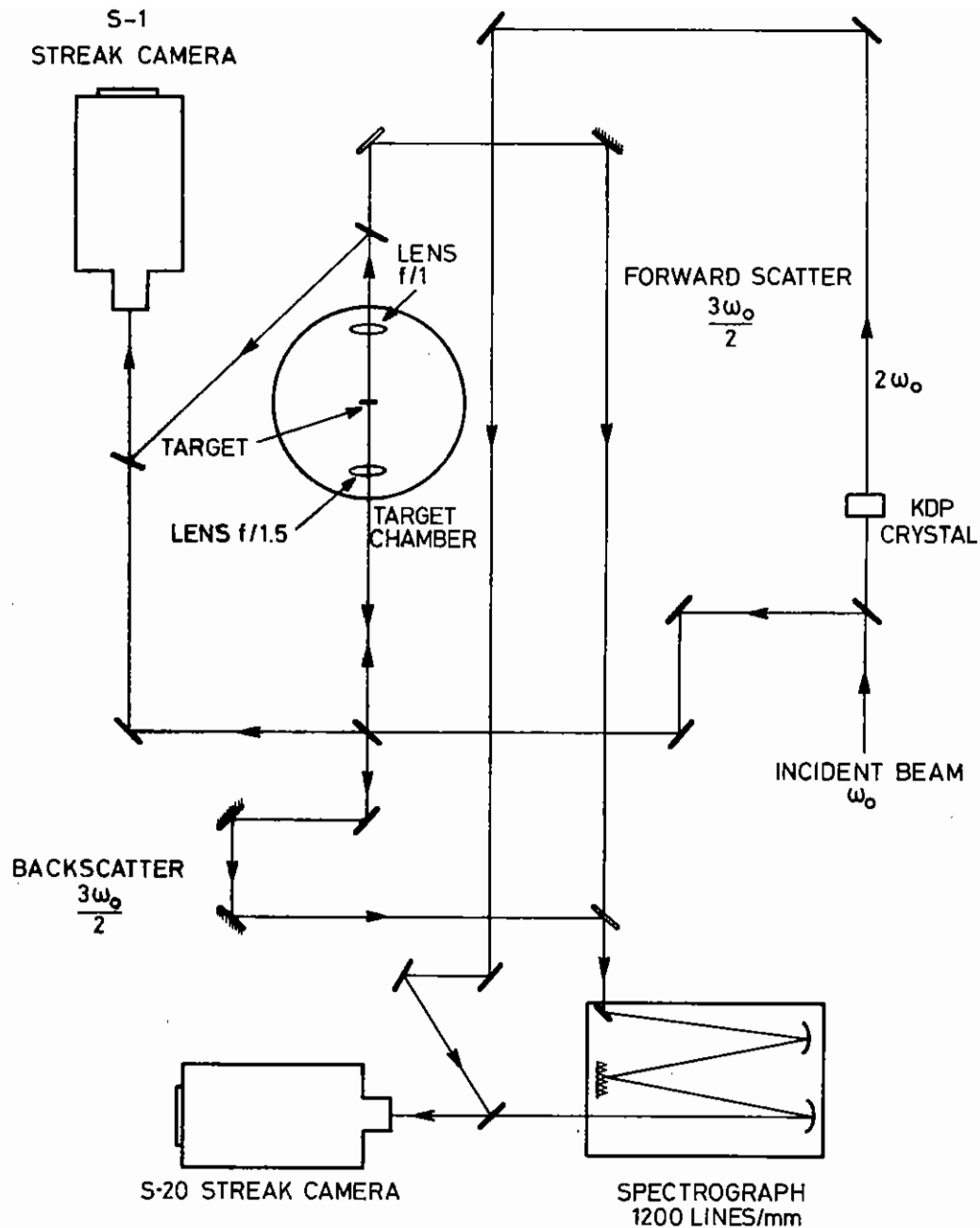


Fig. A1.1

EXPERIMENTAL LAYOUT

$3 \omega_0/2$ harmonic light were spectrally and temporally resolved on an S-20 streak camera with resolutions of $\sim 10 \text{ \AA}$ and $\sim 20 \text{ ps}$ respectively. A portion of the incident light was frequency doubled to provide a time fiducial on the S-20 streak records. The time at which the laser burned through the foil was measured by temporally resolved observations of the incident and transmitted $1.05 \text{ \mu m}$ laser light on an S-1 streak camera.

Figure A1.2 shows the positions of the fiducial, backscattered and forward scattered $3 \omega_0/2$ harmonic spectra on a typical S-20 streak record. Figure A1.3 shows the microdensitometer trace along the time axis of such an example with the corresponding trace of the incident and transmitted light from the S-1 streak camera. The time taken for the plasma to go underdense is estimated from the appearance of the transmitted light, which is normalised to the peak of the reference obtained from a calibration when no foil was present. Although this is only a crude approximation it is sufficient for our purposes since we only want to distinguish between the times when the plasma goes underdense to $9 n_c/4$, n_c and $n_c/4$. Similarly features in the backscattered and forward scattered $3 \omega_0/2$ harmonic spectra are normalised to the peaks of the references obtained for a low energy, frequency doubled shot onto a partially transmitting target.

A1.2.4 Experimental Observations

Various thin foil targets were used - $0.75 \text{ \mu m}$ Aluminium, $0.5 \text{ \mu m}$ and $0.4 \text{ \mu m}$ parylene and $0.35 \text{ \mu m}$ polystyrene - all designed to go underdense during the laser pulse. Figure A1.4 shows microdensitometer traces through the spectra at different times for the S-20 streak record of Figure A1.2 which display the main characteristics of the $3 \omega_0/2$ spectra. Initially, the $3 \omega_0/2$ emission in both forward and backscatter is very broad ($\sim 100 - 150 \text{ \AA}$) and double peaked, pulsating rapidly on a time scale of about 30 ps . Later on the pulsations cease and the spectral width of the $3 \omega_0/2$ emission is appreciably narrower ($\sim 50 \text{ \AA}$). During this later phase of the emission, the backscattered blue wing and the forward scattered red wing become fainter, although the effect is less pronounced in the forward scattered emission. Figures A1.5 and A1.6 show that the timing of this change in the spectra is coincident (within experimental error) with the occurrence of the ω_0

burnthrough, and by implication, with the disappearance of the n_c surface.

A very simple self similar model is used to describe the one dimensional expansion of the foil. The density $\rho(x, t)$ is assumed to be a function only of the self similarity variable $\xi = x/t$ so that the characteristic thickness of the foil, L , is proportional to t and we may write

$$L = C_s t$$

where C_s is the sound speed. The temperature, and thus C_s , varies during the laser pulse and approximately $T \propto \phi^{2/3}$, $C_s \propto W^{1/3}$ so that a first order correction may be made by writing

$$L(t) \propto \int_0^t \phi^{1/3}(t^1) dt^1$$

The peak density in the foil, $\rho_{\max}$, is inversely proportional to L in a one dimensional expansion so

$$\rho_{\max}(t) \propto \frac{1}{\int_0^t \phi^{1/3}(t^1) dt^1}$$

A single measurement of $\rho_{\max} = \rho_c$ at the time of burnthrough then serves to normalise the model for each laser shot. The model is obviously invalid for small times since $\rho_{\max} \rightarrow \infty$ as $t \rightarrow 0$, but it is only used here for densities within a factor of four of ρ_c .

The predicted time of disappearance of the $9 n_c/4$ surface is then close to the appearance of the $3 \omega_0/2$ emission in both back and forward scatter and within the measurement errors it is not possible to say whether the $3 \omega_0/2$ emission begins before or after the $9 n_c/4$ surface has disappeared. The onset of $3 \omega_0/2$ emission is almost certainly determined by the time taken for a large enough density scale length to be set up so that the inhomogeneous threshold for two plasmon decay is exceeded. From the present data it is not possible to determine whether there is any change in the $3 \omega_0/2$ spectrum when the $9 n_c/4$ surface disappears.

TIME AND FREQUENCY RESOLVED $3\omega/2$ EMISSION FROM THIN FOILS

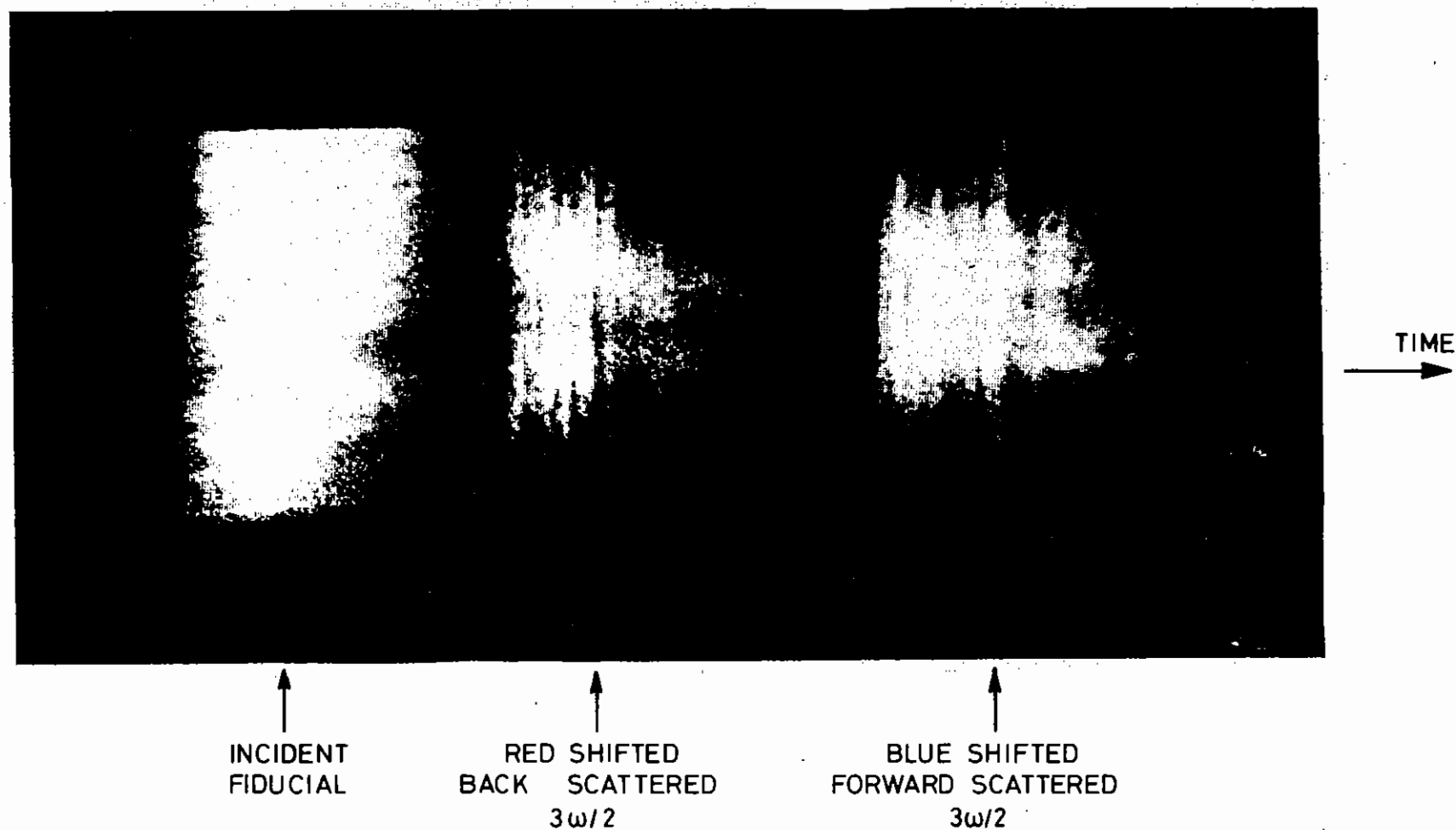
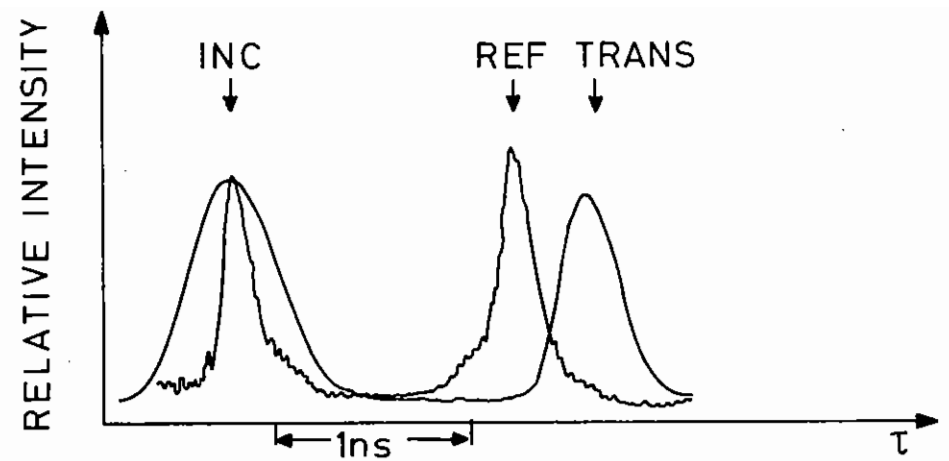
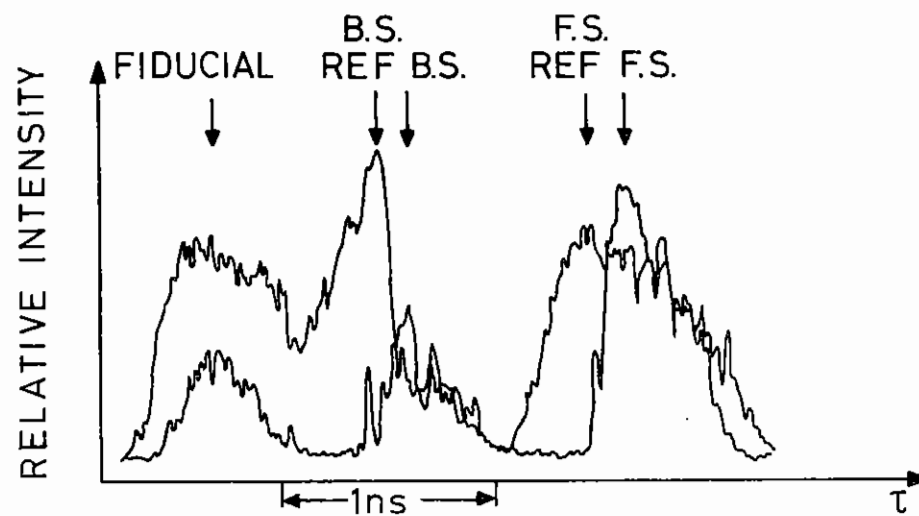


Fig. A1.2



MICRODENSITOMETER TRACE OF ω_0 LIGHT



MICRODENSITOMETER TRACE OF $3\omega_0/2$
HARMONIC EMISSION

Fig. A1.3

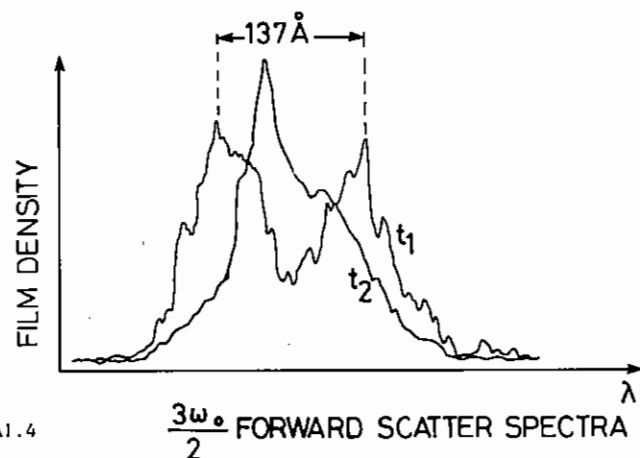
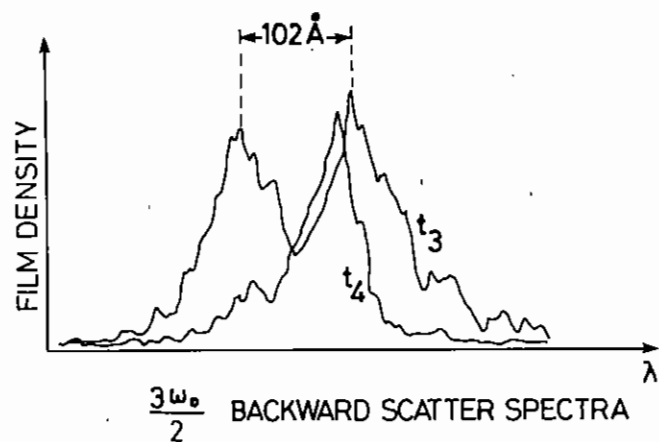
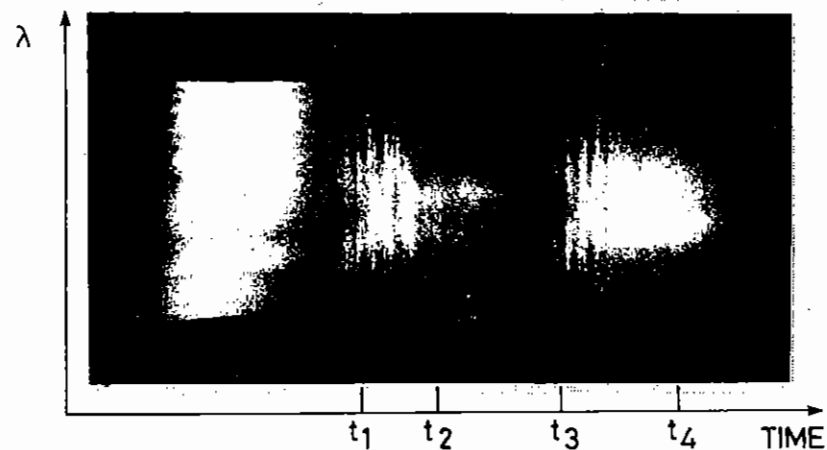
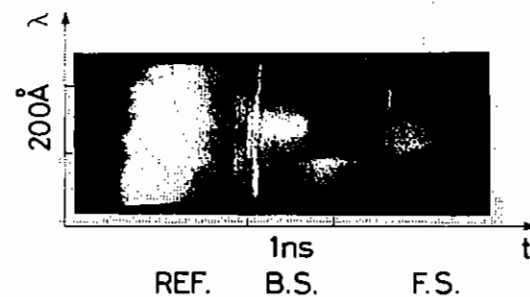


Fig. A1.4

0.35 μm POLYSTYRENE

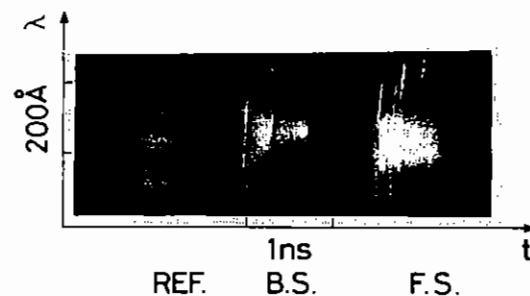


SHOT 4/200583

TIGHT FOCUS

$E_L = 90 \text{ J}$

$\Phi = 2.31 \times 10^{16} \text{ Wcm}^{-2}$

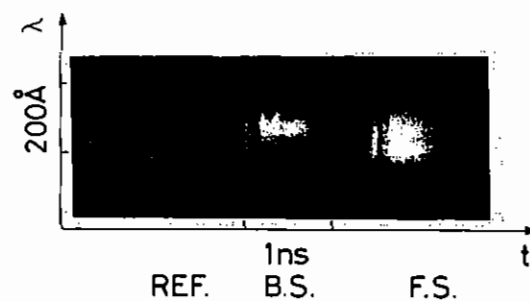


SHOT 8/200583

50 μm FOCUS

$E_L = 97.5 \text{ J}$

$\Phi = 8.68 \times 10^{15} \text{ Wcm}^{-2}$



SHOT 10/200583

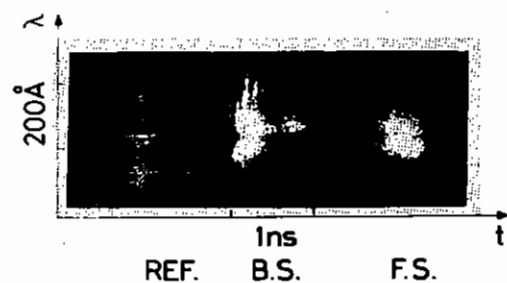
100 μm FOCUS

$E_L = 84 \text{ J}$

$\Phi = 3.46 \times 10^{15} \text{ Wcm}^{-2}$

Fig. A1.5 Variation of signal with intensity, .35 μm target.

0.75 μm ALUMINIUM

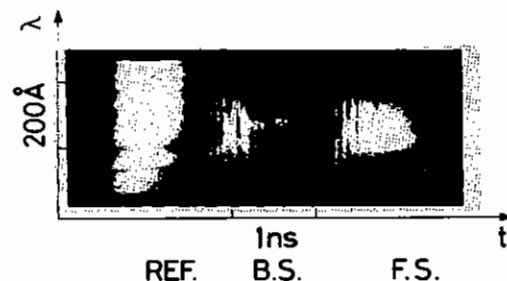


SHOT 11/200583

TIGHT FOCUS

$E_L = 91.5 \text{ J}$

$\Phi = 2.94 \times 10^{16} \text{ Wcm}^{-2}$

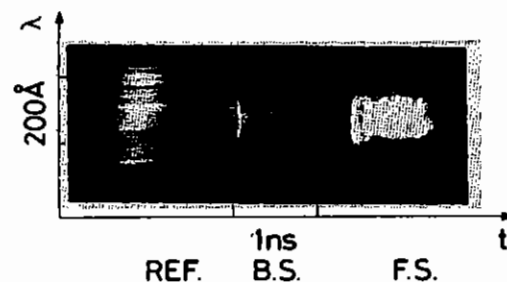


SHOT 3/180583

50 μm FOCUS

$E_L = 69 \text{ J}$

$\Phi = 5.31 \times 10^{15} \text{ Wcm}^{-2}$

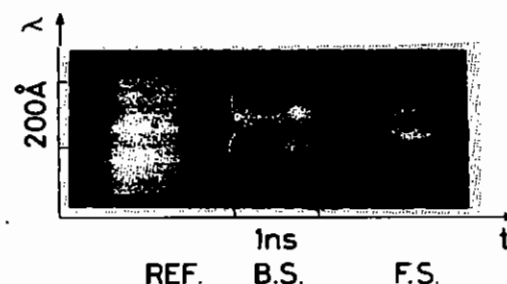


SHOT 10/170583

100 μm FOCUS

$E_L = 64 \text{ J}$

$\Phi = 1.2 \times 10^{15} \text{ Wcm}^{-2}$



SHOT 31/260583

150 μm FOCUS

$E_L = 85 \text{ J}$

$\Phi = 6.04 \times 10^{14} \text{ Wcm}^{-2}$

The same simple model also predicts that the $n_c/4$ surface is present all through the laser pulse and cessation of the $3 \omega_0/2$ emission is due to the pump intensity falling below threshold rather than the plasma density dropping below $n_c/4$.

In summary, the forward and backscatter spectra are double peaked while $9 n_c/4 > n_{\text{max}} > n_c$ and tend to fluctuate on a 30 ps time scale. When $n_c > n_{\text{max}} > n_c/4$ the spectra are also double peaked with a predominance of the red wing in backscatter and the blue wing in forward scatter, but the temporal behaviour has a smooth envelope.

A1.2.5 Theoretical Model

The widely accepted model for $3 \omega_0/2$ generation is the coupling of incident $(\bar{k}_0, \omega_0)$ or reflected $(-\bar{k}_0, \omega_0)$ pump light with $\omega_0/2$ plasma waves produced by two plasmon decay. The ω and $\bar{k}$ matching conditions for this process in a homogeneous plasma have been calculated by Avrov et al (A1.3) and Barr (A1.7). Avrov considers a nonresonant coupling (ie, $k \gg k_0$ where k is the wavenumber of the plasmon), while Barr makes no such assumption. Both models predict a small red-shift from the nominal $3 \omega_0/2$ value for backscattered $3 \omega_0/2$ light and a small blue-shift for forward $3 \omega_0/2$ -emission. Double peak spectra seen in solid target experiments (A1.2, A1.4) can be explained at least qualitatively by the reflection of the forward blue component from the $9 n_c/4$ density (ie, the critical density of $3 \omega_0/2$ light). This is supported by an underdense gas target experiment (A1.8) which shows only a red-shift in backscattered $3 \omega_0/2$ emission. In addition, shifts are larger than predicted by the existing models (A1.3, A1.7), ie the shifts imply rather high electron temperatures at $n_c/4$. This motivated us to study the spectral characteristics of $3 \omega_0/2$ -emission in more detail to interpret experimental observations.

We include for the first time the effects arising from the plasmon propagation and reflection in an inhomogeneous plasma which increase the possible coupling configurations modifying earlier results considerably. In some experimental conditions the pump wave of the two-plasmon process may

Fig. A1.6 Variation of signal with intensity, .75 μm Al target.

travel down the density gradient affecting the spectral behaviour of $3\omega_0/2$ -emission. This type of situation is likely in thin foil targets after burn-through or in underdense gas targets. At high laser intensities also reflected light (from n_c or stimulated Brillouin scattering above $n_c/4$) may exceed the TPD-threshold and act as a pump for the two-plasmon decay.

In the presence of a density gradient the wavenumber k of a propagating plasmon is a sensitive function of the local plasma density particularly in the region of long wavelengths ie, $k^2\lambda_D^2 \ll 1$. From the Bohm-Gross dispersion relation ($\omega^2 = \omega_p^2 + 3k^2v_e^2$, where ω_p is the plasma frequency and v_e the electron thermal velocity) we obtain for the relative change of the wavenumber k

$$\frac{\Delta k}{k} = -\frac{1}{6k^2\lambda_D^2} \frac{\Delta n}{n} \quad (1)$$

where Δn is the change in density. For instance, with $\Delta n/n = 0.05$ and $k\lambda_D = 0.1$ we get $\Delta k/k = -0.83$. On the other hand, for a linear density profile the corresponding distance is $\Delta x = L_n \Delta n/n$ where L_n is the gradient scale length. Thus the maximum distance to the turning point of the plasmon is

$$(\Delta x)_{\max} = 6k^2\lambda_D^2 L_n \quad (2)$$

which is small for $k^2\lambda_D^2 \ll 1$ in a rather steep profile.

This simple argument predicts that the wavelength of a plasmon excited in the two-plasmon process will change considerably during its propagation and may even be reflected from a slightly higher density layer. This means that two-plasmon decay and the following $3\omega_0/2$ generation can take place in slightly separate regions. For instance, a short wavelength $\omega_0/2$ -plasmon ($k \gg k_0$) from TPD is able to couple resonantly with a ω_0 -photon for $3\omega_0/2$ generation at somewhat higher density or after reflection ($k \sim k_0$). Now, the number of different $3\omega_0/2$ coupling configurations is increased and we have four basic cases.

- (i) incident ω_0 -photon and $\omega_0/2$ -plasmon
- (ii) reflected ω_0 -photon and $\omega_0/2$ -plasmon
- (iii) incident ω_0 -photon and reflected $\omega_0/2$ -plasmon
- (iv) reflected ω_0 -photon and reflected $\omega_0/2$ -plasmon

where (i) and (ii) now also include the propagation of the plasmon.

The frequency shift of the $3\omega_0/2$ emission is determined by the matching conditions: $\omega_{3/2} = \omega_0 + \omega_1$ ($3\omega_0/2$ process) and $\omega_0 = \omega_1 + \omega_2$ (TPD) where ω_1 and ω_2 are the frequencies of the TPD-excited plasmons. Using the Bohm-Gross dispersion relation for plasmons and the fact that ω_1 remains constant during the propagation we obtain

$$\omega_{3/2} = \omega_0 + \omega_p \left(1 + \frac{3}{2} k_1^2 \lambda_D^2\right), \quad (3)$$

$$\omega_0 = \omega_p \left(1 + \frac{3}{2} k_1^2 \lambda_D^2\right) + \omega_p \left(1 + \frac{3}{2} k_2^2 \lambda_D^2\right), \quad (4)$$

where ω_p and $k_{1,2}$ are local values corresponding to the density at which two-plasmon instability occurs. Eliminating the plasma frequency ω_p from (3) and (4) we find for the $3/2$ -frequency

$$\omega_{3/2} = \frac{3}{2} \omega_0 + \frac{3}{8} \omega_0 \lambda_D^2 (k_1^2 - k_2^2). \quad (5)$$

From Eq (5) we get the corresponding wavelength shift from the nominal $3/2$ -value

$$\Delta\lambda = -\frac{1}{6} \lambda_D^2 (k_1^2 - k_2^2) \lambda_L \quad (6)$$

where λ_L is the (vacuum) wavelength of the incident laser light. Formally, the shift (A1.6) is the same as given by Barr (A1.7). However the geometrical problem to solve for k_1 and k_2 is quite different when two-plasmon decay and $3\omega_0/2$ generation take place at different densities.

Figure A1.7 shows the k -diagrams for two-plasmon decay and the following

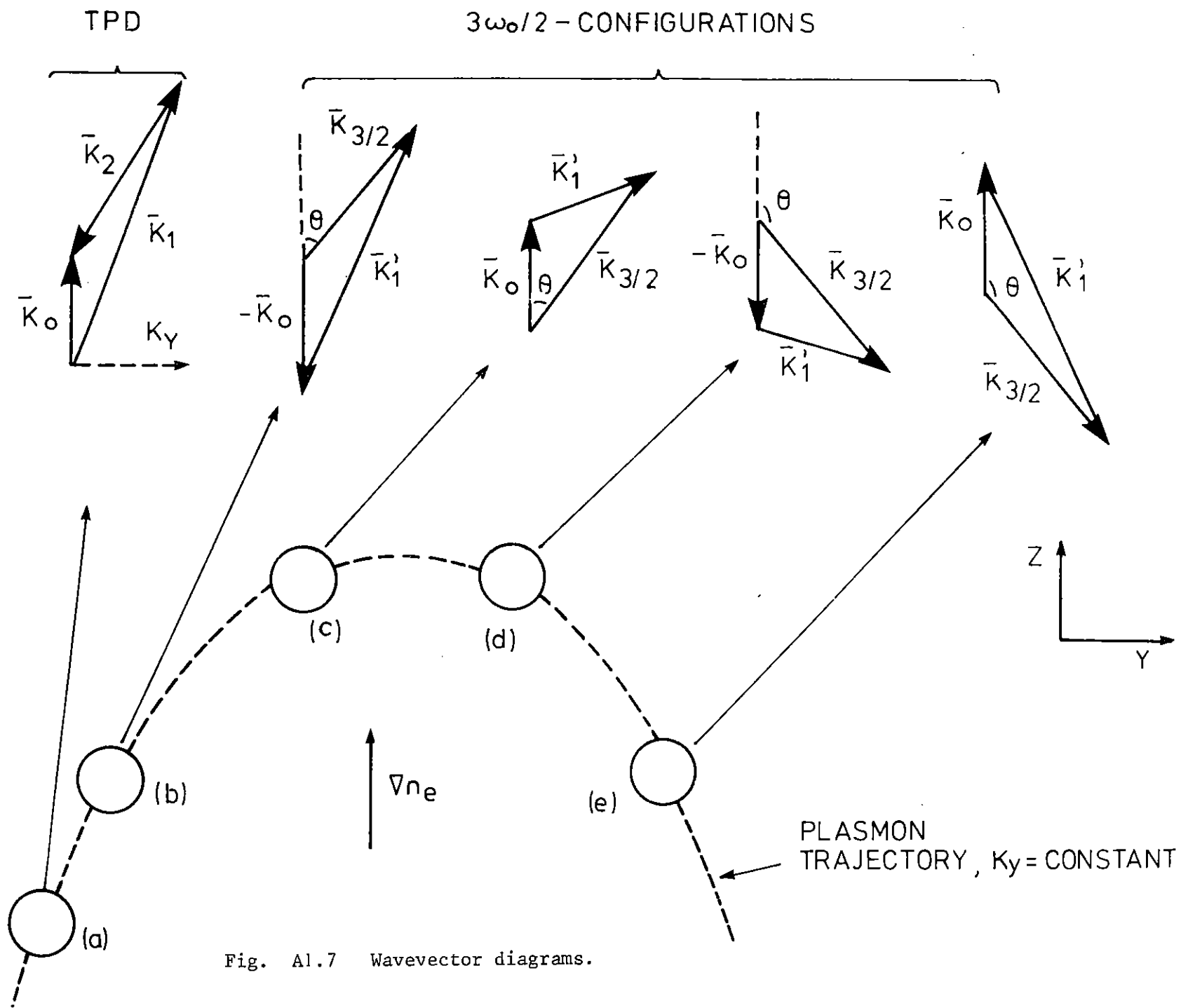


Fig. A1.7 Wavevector diagrams.

$3 \omega_0/2$ coupling geometries with incident or reflected ω_0 -light assuming that TPD operates at a lower density. For simplicity we assume that $\vec{k}_0$ and $\vec{v}_n$ are in the positive z-direction so that the y-component k_y of the plasmon wavevector $\vec{k}_1$ remains constant during the propagation.

We now consider the cases (i) and (iii) given above i.e., $3 \omega_0/2$ generation due to the coupling between the incident photon and plasmon. From Figure A1.7 (a and c or e) we find $k^2 = k_0^2 + k_1^2 - 2k_0 k_1 \cos \theta$ and $k_{12}^2 = k_0^2 + k_1^2 - 2k_0 k_1 \cos \theta$, where θ is the angle of the $3 \omega_0/2$ emission with respect to $\vec{k}_0$ (i.e., small θ corresponds to forward $3 \omega_0/2$ emission and θ near 180° gives backscattered emission). Now we obtain from Eq (6) for the shift

$$\Delta\lambda = \frac{1}{2} \left(\frac{v_e}{c} \right)^2 \lambda_L \left(1 - \frac{2}{\sqrt{3}} (11\alpha^2 - 8\sin^2\theta - 4\sqrt{6} \alpha^2 \cos\theta)^{1/2} \right), \quad (7)$$

where we have used $k_{3/2} = \frac{2}{3} \sqrt{6} k_0$ (k_0 and $k_{3/2}$ are taken constants being much less sensitive to density changes than plasmon wavenumbers). The parameter α in Eq (7) is defined by $\alpha = k_1/k_0$ showing how much larger or smaller the plasmon wavenumber is in the region of TPD excitation as compared to the wavenumber which matches for $3 \omega_0/2$ interaction. The upper limit of α is determined by Landau damping of the plasmon.

From the k-diagrams in Figure A1.7 it can be seen that $3 \omega_0/2$ generation takes place before plasmon reflection for $\theta < \cos^{-1}(k_0/k_{3/2}) \approx 52^\circ$ and after reflection for $\theta > 52^\circ$. Note, however, that the relevant values of the parameter α may be different for these two cases ($\alpha > 1$ for $\theta < 52^\circ$ and $\alpha > 1$ or $\alpha < 1$ for $\theta > 52^\circ$).

The wavelength shift (7) is derived assuming that the incident electromagnetic wave couples to the TPD-plasmon. In a similar way we can analyse the cases (ii) and (iv) when reflected light ($-\vec{k}_0, \omega_0$) couples to $\omega_0/2$ plasmons to generate $3 \omega_0/2$ harmonic emission. The ensuing wavelength shift is

$$\Delta\lambda = \frac{1}{2} \left(\frac{v_e}{c} \right)^2 \lambda_L \left(1 - \frac{2}{\sqrt{3}} (11\alpha^2 - 8\sin^2\theta + 4\sqrt{6} \alpha^2 \cos\theta)^{1/2} \right). \quad (8)$$

Now the angles $\theta < 128^\circ$ correspond to $3 \omega_0/2$ generation before plasmon reflection and $\theta > 128^\circ$ after plasmon reflection (see Figure A1.7). In the limit $\alpha \rightarrow 1$ Eq (8) reduces to $\Delta\lambda = -\frac{1}{2} \left(\frac{v_e}{c} \right)^2 \lambda_L (1 + 4\sqrt{2} \cos\theta)$ which is the same as given by the homogeneous model of Barr (A1.7).

Figure A1.8 shows the wavelength shifts as a function of the $3 \omega_0/2$ emission angle θ for various cases assuming that the incident wave ($\vec{k}, \omega_0$) travels in the direction of the density gradient ($\vec{k} \parallel \vec{v}_n$). The solid curves for $\alpha = 1$ are the shifts without propagation effects of plasmons (A1.7). The dashed part of the curves ($\alpha \neq 1$) represent $3 \omega_0/2$ generation from a propagated plasmon before the turning point of the plasmon and the dash-dotted branches ($\alpha > 1$) are for $3 \omega_0/2$ generation after plasmon reflection. The red-shifts for $\alpha < 1$ in Figure 8 are readily obtained from the coupling between the photon ($\pm \vec{k}_0, \omega_0$) and the "red" plasmon ($\vec{k}_2, \omega_2$) which may not be excited at a lower density than that at which $3 \omega_0/2$ emission originates.

A new feature is that a relatively large blue-shift in backscattered $3 \omega_0/2$ emission can exist simply due to the reflection of the "blue" plasmon ($\vec{k}_1, \omega_1$). The shifts are also modified by the propagation effects of plasmons such that blue-shifts tend to enhance and red-shifts to reduce, thus reducing the asymmetry of red and blue shifts given by the homogeneous plasma model.

In some experimental conditions the incident pump wave may propagate down the density gradient. This is the case for underdense targets when the peak density is less than the critical density so that there are quarter critical layers at both sides of the peak density. Density ripples in underdense plasma may also lead to such a situation. At high laser irradiance the reflected light may become intense enough to exceed the threshold of two-plasmon decay. In this case also the pump wave (now reflected light) travels down the density gradient.

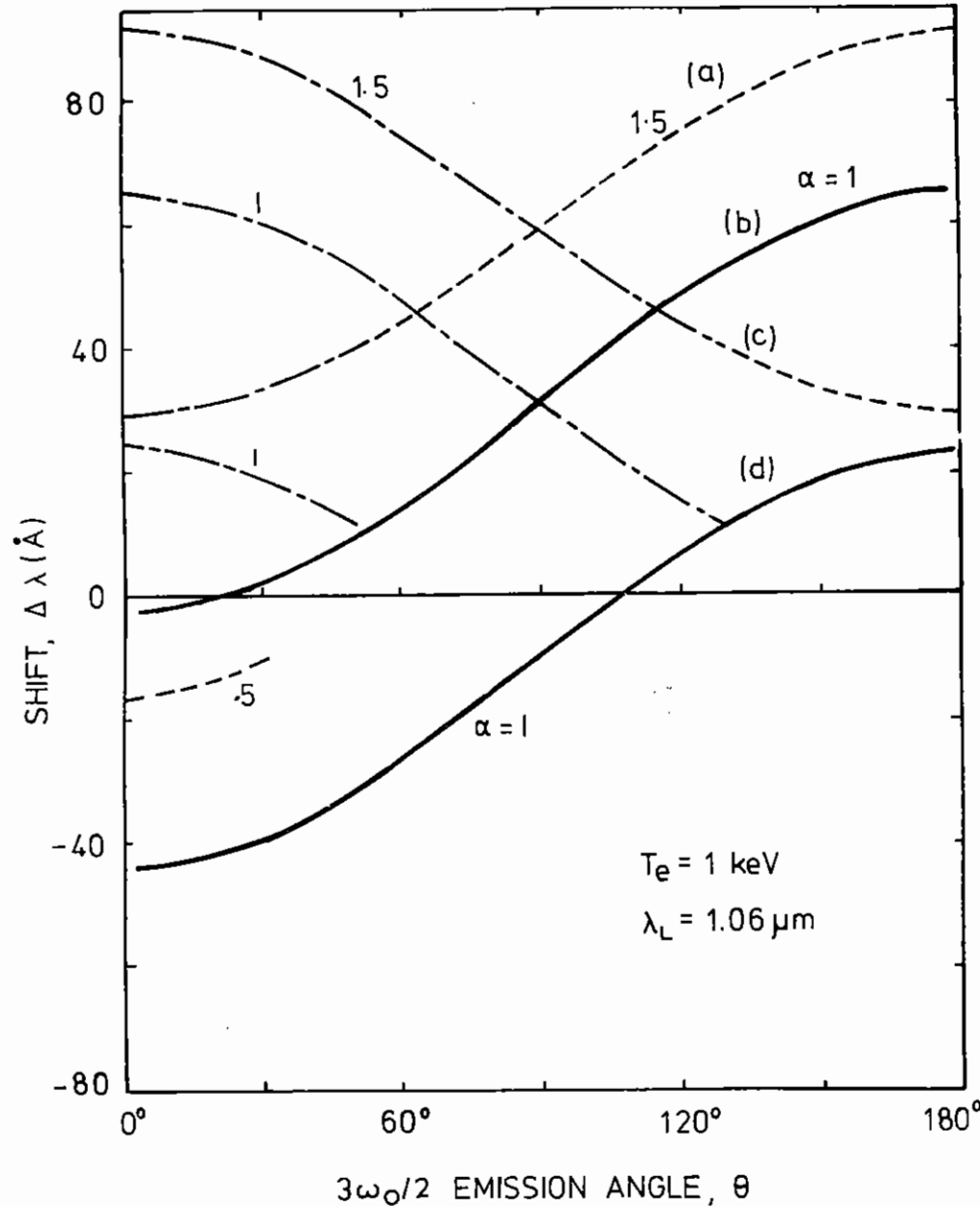


Fig. A1.8 Wavelength shift vs emission angle k_0 parallel to ∇n .

In this type of condition the "red" plasmon ($\bar{k}_2, \omega_2$) propagates to higher densities and can be reflected. The ensuing wavelength shift due to coupling to the plasmon ($\bar{k}_2, \omega_2$) is

$$\Delta\lambda = \frac{1}{2} \left(\frac{v_e}{c} \right)^2 \lambda_L \left(1 - \frac{2}{\sqrt{3}} (11\alpha^2 - 8\sin^2\theta \pm 4/6 \alpha^2 \cos\theta)^{1/2} \right), \quad (9)$$

where the "-" sign refers to the $3\omega_0/2$ -coupling between the pump wave ($\bar{k}_0, \omega_0$) and plasmon and the "+" sign for the photon ($-\bar{k}_0, \omega_0$).

Figure A1.9 illustrates the shifts versus $3\omega_0/2$ emission angle θ for the case of this reversed density gradient ($\bar{k}_0 \parallel -\nabla n$). The solid curves ($\alpha = 1$) are the same as in Figure A1.8, the dashed part of the curves ($\alpha > 1$) are from the $3\omega_0/2$ emission before plasmon reflection and the dash-dotted branches ($\alpha > 1$) are from the coupling to reflected plasmons. As expected the red-shifts are now increased and blue-shifts decreased. Also a red-shifted component in forward $3\omega_0/2$ emission appears due to reflection of the "red" plasmon which was not present in the previous case of normal incidence (cf Figure A1.8).

Finally we consider the case in which the density gradient is perpendicular to the incident wavevector ($\bar{k}_0 \perp \nabla n$). This is likely in filaments or in the presence of a strong self-focusing when the two-plasmon decay occurs on the wall of the plasma channel. The plasmons propagate such that the k_z -component of the plasmon wavevector (ie, component parallel with $\bar{k}_0$ and perpendicular to ∇n) remains constant. Simple geometry gives $k^2 - k_z^2 = k_0^2 + 2k_0 k_z$ which is independent of the α -parameter. Thus the shift (6) does not depend on the plasmon propagation or reflection effects and the shifts are given by the solid curves in Figures A1.8 and A1.9.

To summarise the features of the model:

(a) In backscatter, the blue peak requires that either the $3\omega_0/2$ photon be reflected at $9n_c/4$ or that the forward propagating plasmon be reflected close to $n_c/4$.

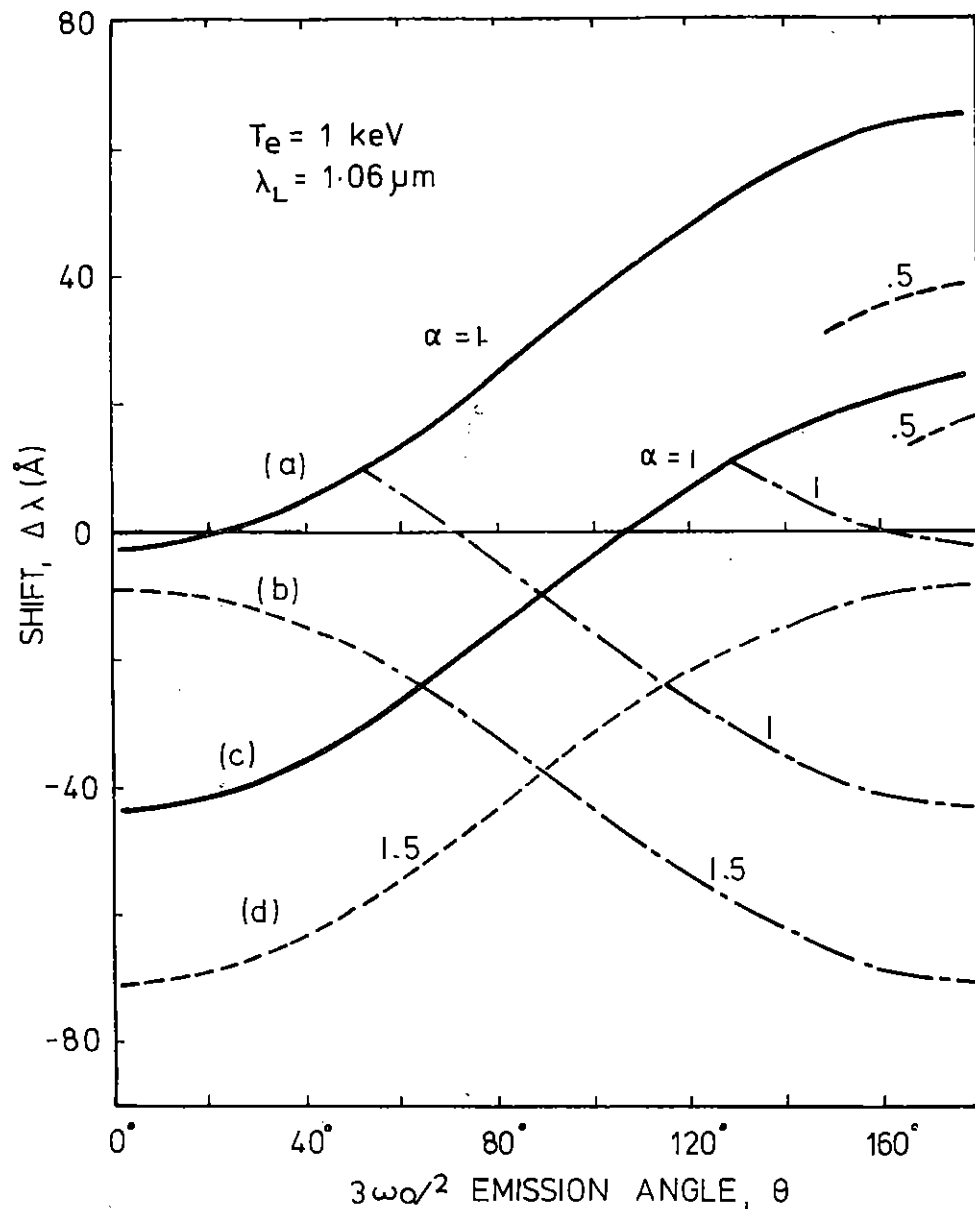


Fig. A1.9 Wavelength shift vs emission angle k_0 parallel to $-\nabla n$.

(b) In forward scatter, with the pump wave travelling up the density gradient, only the blue shifted peak is generated. The red peak requires that the pump wave for the two plasmon decay process be travelling down the density gradient. Experimentally this may occur if the reflected laser light is sufficiently intense to exceed the two plasmon decay threshold, or on the "back" side of a thin foil after burnthrough. In this latter case the "red shifted" plasma wave must reflect before combining with a pump photon.

(c) The effect of plasma wave propagation in an inhomogeneous plasma is to make the red and blue shifts more nearly equal than in the homogeneous plasma model of Barr (Ref 7). The uncertainty in the value of α in equations (8) and (9) reduces the value of the $3\omega_0/2$ spectrum as a temperature diagnostic.

A1.2.6 Conclusions

In solid target experiments ((A1.3), (A1.4), (A1.5)) the multiplicity of possible processes makes it difficult to establish the exact coupling mechanisms leading to $3\omega_0/2$ photon or the $\omega_0/2$ plasmon.

In gas target experiments (A1.2) the reflection of the $3\omega_0/2$ photon is impossible since there is no $9n_c/4$ surface and reflection of the $\omega_0/2$ plasmon is small if the density gradients are very long. Experimentally only the red peak is seen in backscatter in gas targets.

The thin foil experiment reported here shows both red and blue wings in backscatter after the peak density in the foil has dropped below $9n_c/4$, showing clearly that propagation and reflection of the $\omega_0/2$ plasma waves is occurring. The blue wing generally decreases in intensity at late times when the foil develops longer density gradients. As the separation of the points of generation and reflection of the TPD plasmons increase, their absorption (both collisional and Landau) along the increased path length is greater and consequently the blue wing becomes fainter.

The observation of the red wing in forward scatter requires that in some regions the pump wave is travelling down the density gradient. There are three circumstances in which this could arise.

- (a) The laser light reflected from critical density is sufficiently intense to drive TPD.
- (b) After the foil has burned through the pumplight penetrates through to the second $n_c/4$ layer on the rear surface.
- (c) The two plasmon decay on the front of the foil may saturate by forming a local density depression with larger densities in both the forward and backward directions.

Processes (a) and (b) tend to be exclusive and one would expect a change in the appearance of the forward scattered red wing around the time of burnthrough. There is a pronounced change in the $3\omega_0/2$ spectrum around the time of burnthrough but since it affects forward and backscattered red and blue peaks it is not a conclusive indicator of the switch from process (a) to process (b). Profile modification around $n_c/4$ due to the ponderomotive pressure of the electrostatic decay waves must occur to some extent, and only a small density depression, $\Delta n/n \approx 0.1$ is required to reflect the plasma waves. Process (c) above must be regarded as a likely explanation of the forward scattered red wing but again the results are not definitive.

There remains the dramatic change in the $3\omega_0/2$ spectrum occurring around the time of the disappearance of the critical density surface. While the n_c surface is present the spectrum is very broad and pulsates rapidly in time, after burnthrough the spectrum is narrower, but still double peaked, and varies smoothly in time. These rapid pulsations have been observed previously in $2\omega_0$ (A1.9), $3\omega_0/2$ (A1.4), ω_0 (A1.10) and $\omega_0/2$ (A1.11), from solid targets and explanations involving filamentation, unstable hydrodynamic flow at critical density or bursts of hot electrons have been suggested. The occurrence of filamentation is normally associated with long scalelengths in the underdense plasma whereas these oscillations are observed predominantly at early times while the plasma scale lengths are

short and we regard this explanation as unlikely. Since the same behaviour is observed in processes originating at n_c (1e ω_0 and $2\omega_0$ emission) and at $n_c/4$ ($3\omega_0/2$ and $\omega_0/2$ emission) the pulsations are unlikely to arise out of any mechanism peculiar to the two plasmon decay but are more likely due to a longer range effect. Such an effect can be provided by an unstable fluid flow across the ponderomotive pressure jump at critical density. This has been seen in fluid simulations (A1.12) and predicted by a steady flow model (A1.13). This instability may either affect the ablation flow and thus propagate its influence out into the underdense plasma, or alternatively the changing density profile will change the resonance absorption conditions and may introduced pulsations into the source of hot electrons. The hot electron density in the coronal plasma is essentially uniform because of their long mean free path, and the hot electron density at quarter critical will respond rapidly to changes in the source term. Since the Bohm Gross dispersion relation is essentially an electron pressure correction a changing hot electron population will alter the k and ω of the most unstable two plasmon decay waves. A flow instability at critical density or a pulsating source of hot electrons can thus influence processes occurring at n_c and $n_c/4$ simultaneously. This can be either through the fluid dynamics of the expansion flow or through the influence of hot electrons on the plasma dispersion relation. Removal of the critical density surface removes both the ponderomotive pressure jump and the source of hot electrons and thus also removes the cause of the pulsating harmonic emissions. On the basis of the present data those two mechanisms cannot be distinguished but a pulsating source of hot electrons could be detected through the hard X-ray Bremsstrahlung with a suitable X-ray streak camera.

Quite separately from any question of pulsations, the hot electrons generated at critical density will make a substantial contribution to the width of the $3\omega_0/2$ emission, which is predicted to be proportional to T_e , or more generally to P_e if the electron distribution function is not Maxwellian.

Experimental measurements (A1.14, A1.15) of the fraction of incident laser energy going into hot electrons η_H , give $0.05 < \eta_H < 0.1$ for short pulse experiments at $\phi = 2 \times 10^{15} \text{ W cm}^{-2}$ and $\lambda = 1 \text{ }\mu\text{m}$. The density of the hot

electron population n_H is estimated by equating the energy carried away by the hot electrons in the free streaming limit with the absorbed power into hot electrons.

$$n_H \phi = f n_H kT_H v_H$$

where f is a flux limiting factor which for hot electrons in a cold background we take to be $f = 0.6$. (Any reduction in f would lead to an increase in the hot electron density n_H .) Since the hot electron mean free path is large compared with the separation of n_c and $n_c/4$ the hot electron density in this region will be approximately uniform.

Typically at $\phi = 2 \times 10^{15} \text{ W cm}^{-2}$; $T_H = 12 \text{ keV}$ (A1.15) and the above equation gives $n_H = 2.5 \times 10^{19} \text{ cm}^{-3}$, compared with $n_c/4 = 2.5 \times 10^{20} \text{ cm}^{-3}$. T_H is typically ten times the background thermal electron temperature T_c so in the presence of hot electrons the electron pressure

$$P_e = n_c (0.9 kT_c + 0.1 kT_H)/4$$

is about twice the pressure in the absence of hot electrons

$$P_e = n_c kT_e/4$$

The reduction of the spectral width of $3 \omega_0/2$ emission by about a factor of two when the plasma becomes underdense is thus explained by the change in the total electron pressure when the hot electrons are removed.

One further effect of the hot electron production is to drive a large cold electron return current to maintain quasi-neutrality. The return current may drive ion acoustic turbulence and the ion waves ($\omega_{ia} < \omega_{pe}$) act in turn to broaden the spectrum of electron plasma waves. This effect is probably significant in the broadening of $2 \omega_0$ emission (A1.9) but is probably not a major contributor to the much larger width of the $3 \omega_0/2$ emission.

In summary we believe that the change in width of $3 \omega_0/2$ emission coincident with the disappearance of the n_c surface is explained by the contribution of

hot electrons to the plasma dispersion relation. Pulsations in $2 \omega_0$ and $3 \omega_0/2$ emission in the presence of a steep density profile at n_c are probably linked and caused by rapid changes in the density profile at n_c .

A1.3 Raman Scattering from Pre-formed Plasmas

V Aboites, E McGoldrick, S M L Sim (Essex), D Bassett, A J Cole and C J Hooker (RAL)

In laser produced plasmas two processes are thought to give rise to plasma emission in the region of $\omega_0/2$. These are the Two Plasmon Decay Instability (TPDI) (A1.5) and Stimulated Raman Scattering (SRS) (A1.6) both of which occur near the quarter critical density region. In addition SRS can occur at much lower densities (A1.16). Both instabilities may generate very fast electrons which are detrimental to laser fusion.

Generally the TPDI dominates, except for very high irradiances and in long scalelength plasmas so earlier attempts at Rutherford to observe SRS have failed. To overcome this problem large scalelength plasma were produced by using pre-formed plasma. Here thin foils were exploded with a pre-pulse, and the main pulse was delayed until the maximum density of the plasma had fallen below critical density (measured using burn through techniques). The experiment then spectrally and temporally resolved emission in the region of $\omega_0/2$.

Foil targets of $0.7 \mu\text{m}$ to $0.3 \mu\text{m}$ thick parylene were irradiated with two pulses of $0.53 \mu\text{m}$ laser light travelling in opposite directions. The pre-pulse (15 - 25 J, $\sim 420 \text{ ps}$ duration) formed the plasma and a nanosecond later the second pulse (40 - 60 J, $\sim 560 \text{ ps}$ duration) with irradiances between 10^{13} and $10^{15} \text{ W cm}^{-2}$ hit the plasma (Figure A1.10).

Emission from the plasma was collected by the focusing lens of the second pulse ($f/1.5$) and directed onto the slits of a grating spectrograph (300 $1/\text{mm}$ grating). The output from the spectrograph was then coupled to the S1 streak camera to give a spectral window of $5500 \text{ \AA}$, centred at $9000 \text{ \AA}$. The camera was operated at 600 ps cm^{-1} streak speed giving spectral and temporal

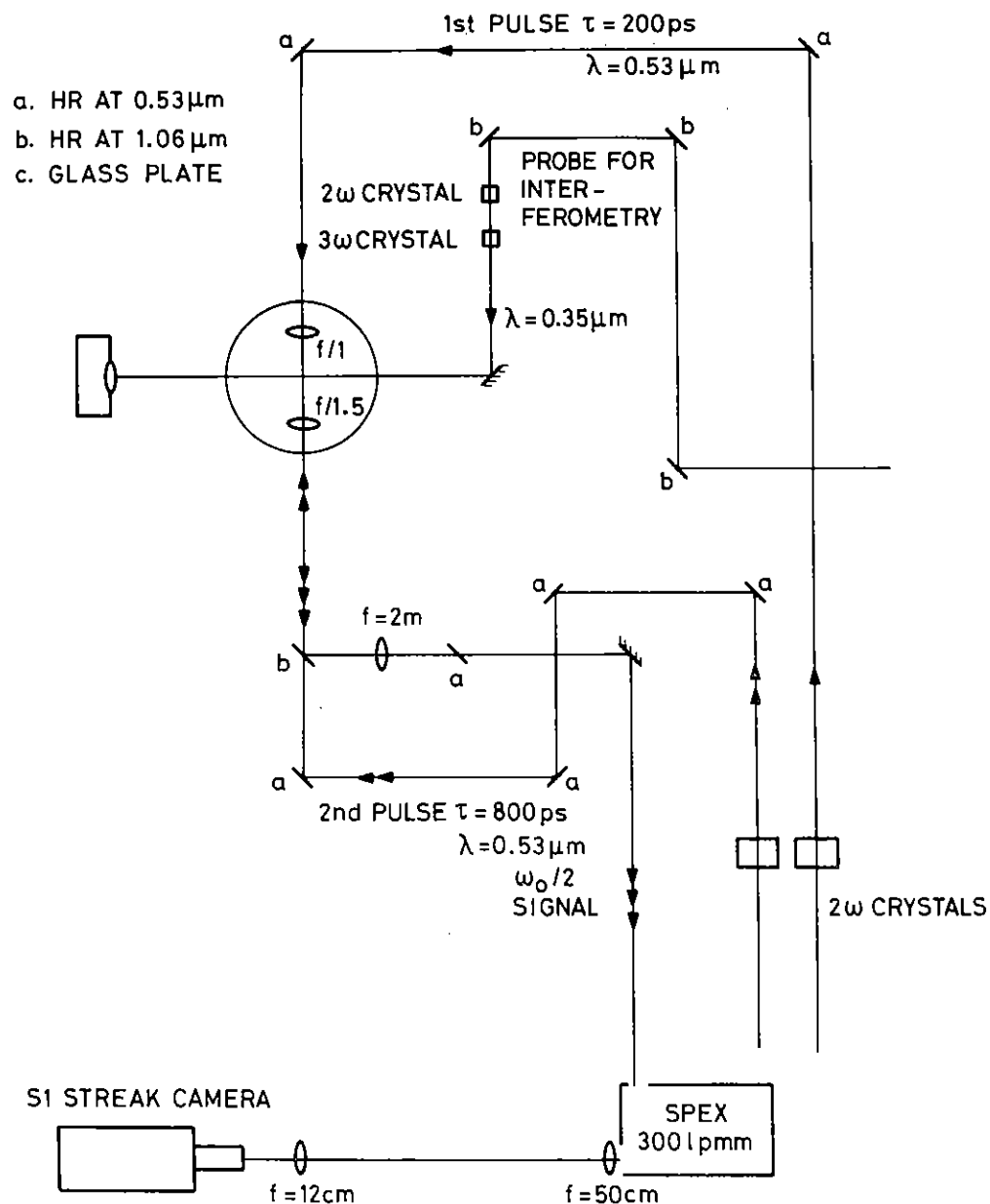


Fig. A1.10 Experimental layout for SRS.

resolutions of $35\text{ \AA}$ and 120 ps respectively.

Other diagnostics used were an X-ray diode array to measure hard X-rays and interferometry (using a $0.35\mu\text{m}$ probe) to measure scalelengths in the plasma.

Typically the streak record is expected to consist of three signals; unconverted $1.05\mu\text{m}$ light (acting as a spectral and temporal marker), forward scattered emission produced by the pre-pulse, and back-scattered emission from the main pulse.

Figure A1.11 shows a typical streak record for a $0.7\mu\text{m}$ foil. It can be seen that the forward scattered emission is relatively narrow, double peaked and centred round the nominal $\omega_0/2$ frequency. The backscattered emission is very different, being very broad (FWHM $\sim 3000\text{ \AA}$), single peaked and shifted to shorter wavelengths with the peak at $\sim 8800\text{ \AA}$. This emission is almost certainly due to convective SRS occurring in the low density regions of the plasma upto quarter critical density (A1.17).

The first signal is similar to that observed in earlier $\omega_0/2$ experiments (A1.18) and is attributed to scattering of the pump wave from plasmons created at quarter critical density by the TPDI in a process similar to that which produces $3\omega_0/2$ emission (A1.19).

For the $0.3\mu\text{m}$ thick foil a typical streak record is shown in Figure A1.12. Here little forward scattered emission is observed probably because the quarter critical density exists for a much shorter time and so less TPDI can occur there. The backscattered emission is similar to that seen for the $0.7\mu\text{m}$ foil and is due to SRS. Here the FWHM is $\sim 1500\text{ \AA}$, and peaks at $\sim 9270\text{ \AA}$.

The temporal evolution of the forward and backscattered signals was analysed for the $0.7\mu\text{m}$ foil shot shown earlier. It was found that the peak of the SRS emission occurred at the peak of the main pulse as expected, and that the peak of the TPDI emission occurred after the peak of the pre-pulse. This is because there would have to be time for the scalelengths at quarter critical

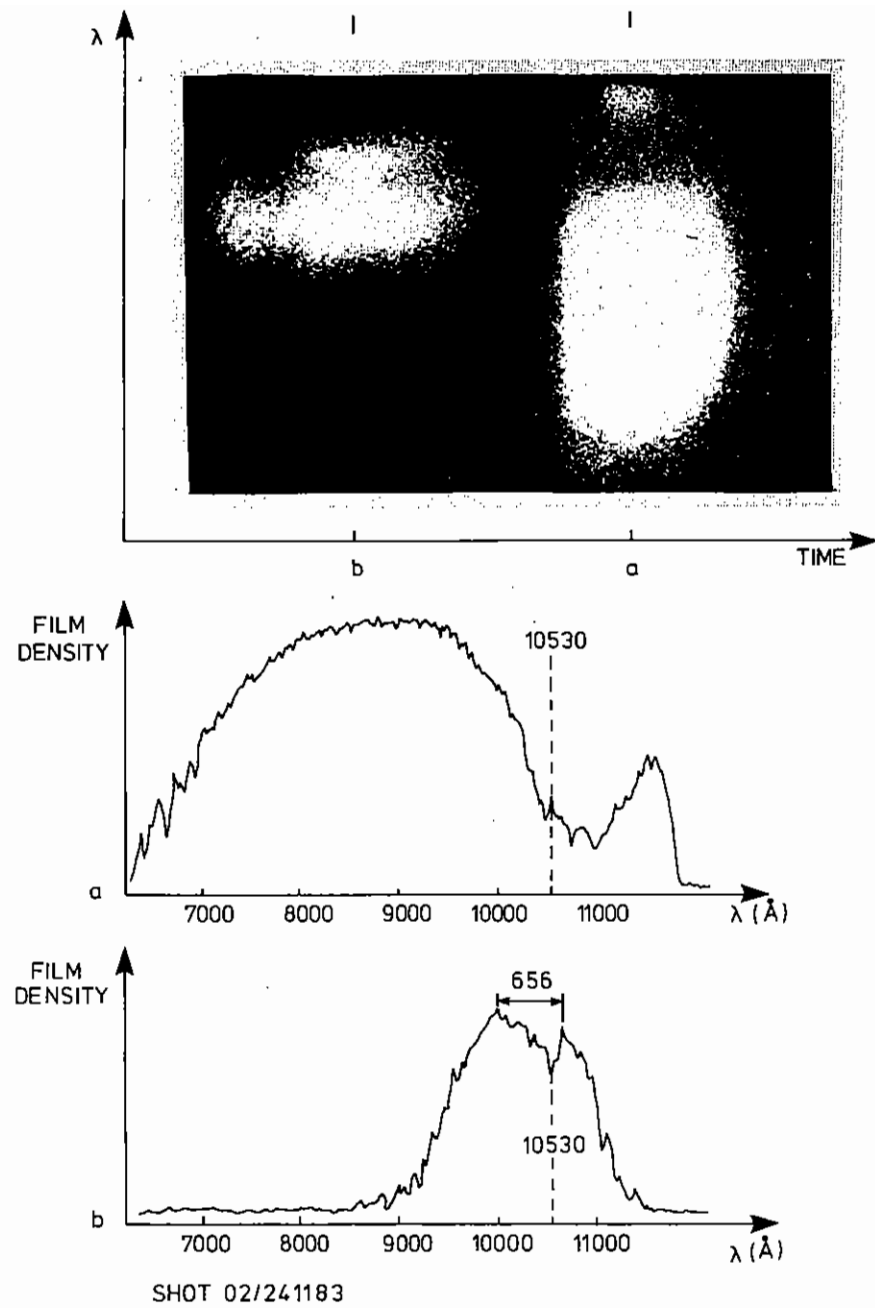


Fig. A1.11 Results obtained for a $0.7\mu\text{m}$ foil
 a) microdensitometer trace of the backscattered emission
 b) microdensitometer trace of the forward scattered emission.

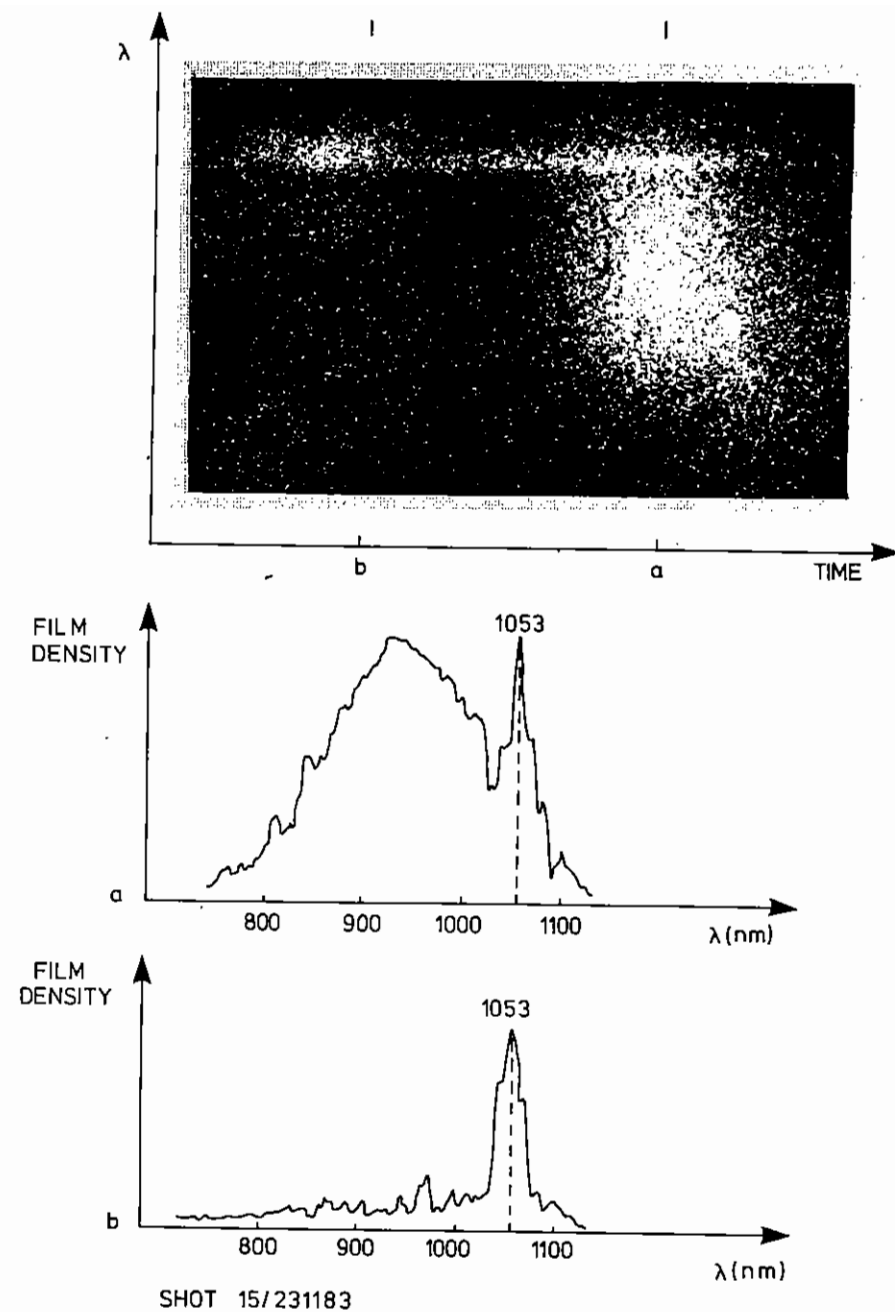


Fig. A1.12 Results obtained for a $0.3\mu\text{m}$ foil
 a) microdensitometer trace of the backscattered emission
 b) microdensitometer trace of the forward scattered emission.

density to relax enough for the instability threshold to be reached.

Unfortunately little information could be obtained from the interferometry. Analysis has shown that the probe was unable to reach sufficiently high densities to be of use. The X-ray diode measurements gave some indication that harder X-rays were observed when SRS was observed, but too few data shots were obtained, and the signal levels were too low for any more definite conclusions to be reached.

Summing up, this experiment has shown that, by using this technique to produce long scalelength plasmas, conditions can be created in which SRS can be made to dominate over the TPDI without the need for very high irradiances. Further work is necessary however, before a full understanding of these processes is reached.

Al.4 Thomson Scattering from Two Plasmon Decay Waves in a 1.05 μm Laser Produced Plasma

V Aboites, S M L Sim (Essex), H A Baldis (NRC, Ottawa) and R G Evans (RAL)

Scattered light at $3\omega_0/2$ harmonic can be generated by the coalescence of a pump or reflected light wave with $\omega_0/2$ electron plasma waves produced by the two plasmon decay instability at the $n_c/4$ layer. Although the growth rate for the two plasmon decay instability is highest for $k_p \gg k_o$, where the subscripts p and o denote the decay and pump waves respectively, the matching conditions give the maximum growth rate for $3\omega_0/2$ harmonic generation for $k_p \sim k_o$.

Previous observations (Al.4, Al.20) of the $3\omega_0/2$ spectra have shown that it occurs in short bursts of emission on a ~ 30 ps timescale. Similar pulsations have been observed in time resolved Thomson scattering studies of the k-spectra of plasma waves driven by the two plasmon decay instability in a CO_2 laser produced plasma (Al.21). This experiment was designed to measure the k-spectra of the two plasmon decay waves involved in $3\omega_0/2$ harmonic generation by correlating the temporal behaviour of simultaneously observed $3\omega_0/2$ and Thomson scattered signals.

Al.18

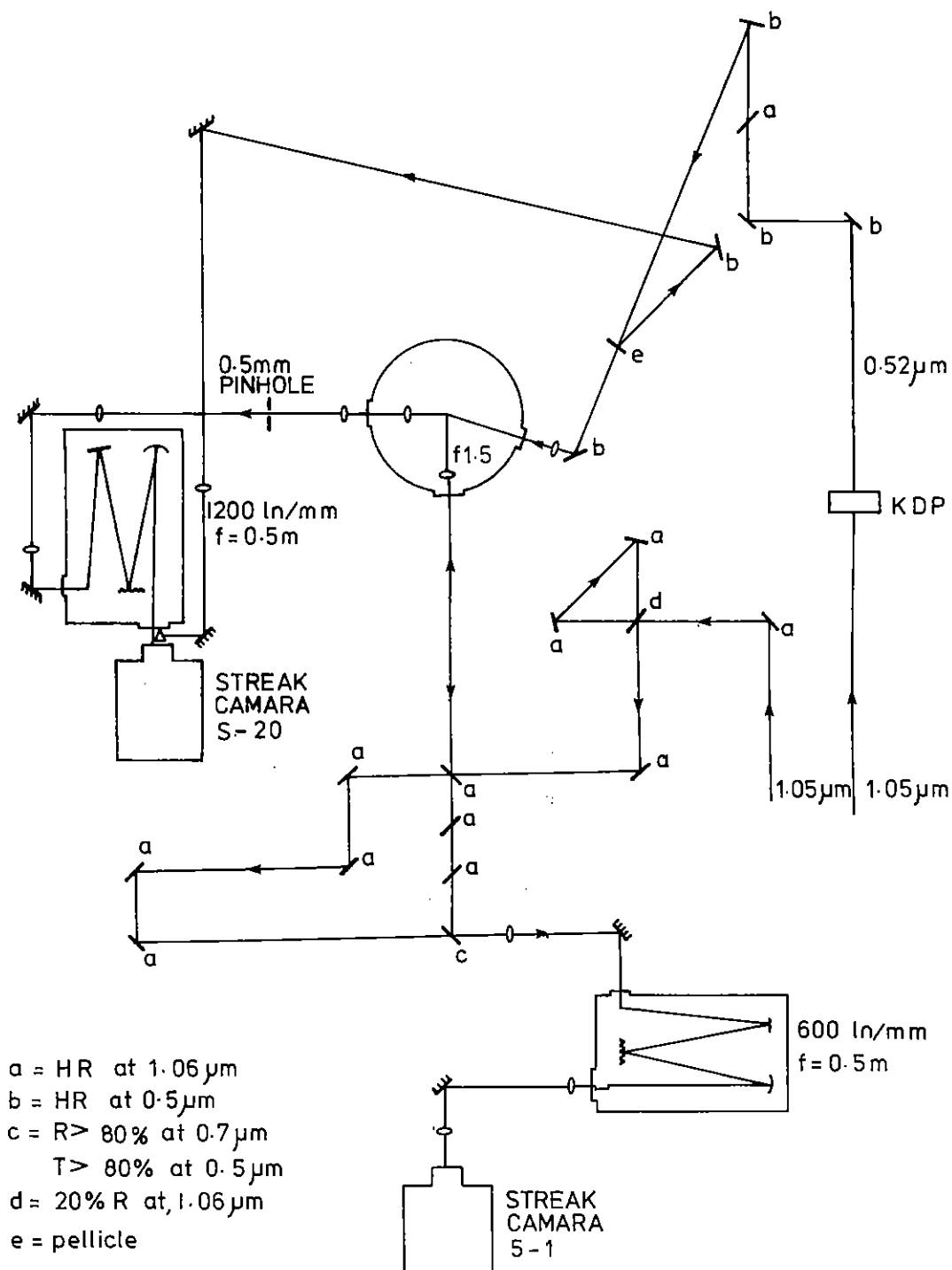


Fig. Al.13 Experimental layout for Thomson scattering.

The experimental layout is shown in Figure A1.12. The backscattered $3 \omega_0/2$ harmonic spectra, produced from 5 μm thick Al foil targets irradiated by 1.05 μm , 1.8 ns laser pulses, were recorded on an S-1 streak camera. An S-20 streak camera was used to record the Thomson scattered signals from an 0.5 μm , 1.5 ns probe beam which was synchronous with the main 1.05 μm beam. In order to reduce the level of the $5 \omega_0/2$ harmonic signal an etalon was introduced into the 1.05 μm beam to generate a pre-pulse containing $\sim 20\%$ of the incident energy. The separation between the pre-pulse and the main pulse was arranged so that the peak plasma density n_e was $n_c < n_e < n_c/4$ at the time of arrival of the main pulse.

However, despite these precautions, the levels of $5 \omega_0/2$ signals and refracted green light were still comparable with the Thomson scattered signals making it difficult to identify a true "Thomson scattered" signal.

A1.5 Detailed Measurement and Analysis of Near Classical Thermal Transport in 1 μm Irradiated Laser Plasmas

J D Kilkenny, O Landen, S D Tabatabaei, J Wark (IC), O Willi, A Hauer (LANL, USA) and C J Hooker (RAL)

The transport of energy from the region of the absorption to denser portions of laser produced plasmas has remained a topic of active investigation and controversy for nearly ten years (A1.22, A1.23, A1.24). Early work largely involving flat targets seemed to indicate a severe limitation of the heat flux relative to classical predictions (A1.25). More recent work involving spherical targets has tended to yield transport results closer to classical (A1.26, A1.27, A1.28). In the present experiment we have studied transport and plasma evolution in spherical targets using a wide variety of diagnostics. This wide range of diagnostics yields a more comprehensive picture of the plasma conditions than has been presented previously. For example, the combination of time resolved X-ray line spectroscopy (as a diagnostic of thermal transport) and interferometric probing (as a diagnostic of plasma expansion) places important constraints on the theoretical modelling.

The basic target configuration used in most of the work is shown in Figure A1.14. The solid glass ball prevented any implosion of the target (and consequent inward movement of the critical surface) during the shot. In most cases the targets were uniformly illuminated with the six-beam laser facility at the Rutherford Appleton Laboratory. The six beams are focused by f/1 lenses with optical axes normal to the faces of a cube. Incident irradiance levels ranged from 3 to $8 \times 10^{14} \text{ W/cm}^2$. The laser pulse had an approximately Gaussian temporal profile with a 0.83 ns halfwidth.

The primary diagnostic of thermal transport was time resolved X-ray line spectroscopy. X-ray line emission was observed primarily from silicon in the glass centre and the thin aluminium layer. A few lines were also observed from sodium, calcium and magnesium impurities in the glass. The experimental configuration used for time resolved spectroscopy is shown in Figure A1.15. A dispersing crystal is located about midway between the target and the streak camera slit. Also shown in Figure A1.15 is the method used for obtaining a time fiducial mark indicating when light is first incident on the target.

Another important diagnostic used in these experiments was interferometric probing. A 40 ps, 3ω probe was used to measure the density profile in the expanding plasma. In addition to giving details on the expansion profile, this serves to show that quite uniform blowoff conditions have been achieved in the case of uniform illumination.

Other diagnostics that were employed are:

- (a) Ion calorimetric measurement of absorbed energy. For uniform illumination, the absorbed fraction was about 27%.
- (b) Faraday cup measurement of ion blow-off.
- (c) Faraday rotation measurement of magnetic fields in the cases when nonuniform illumination was employed.
- (d) Multiple channel pinhole X-ray imaging.

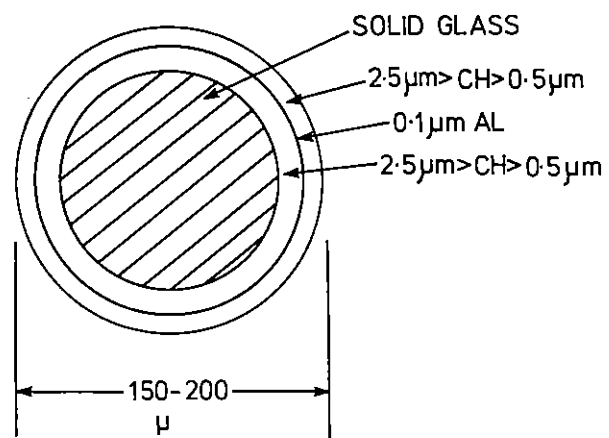


Fig. A1.14 Basic target configuration used in electron transport experiments.

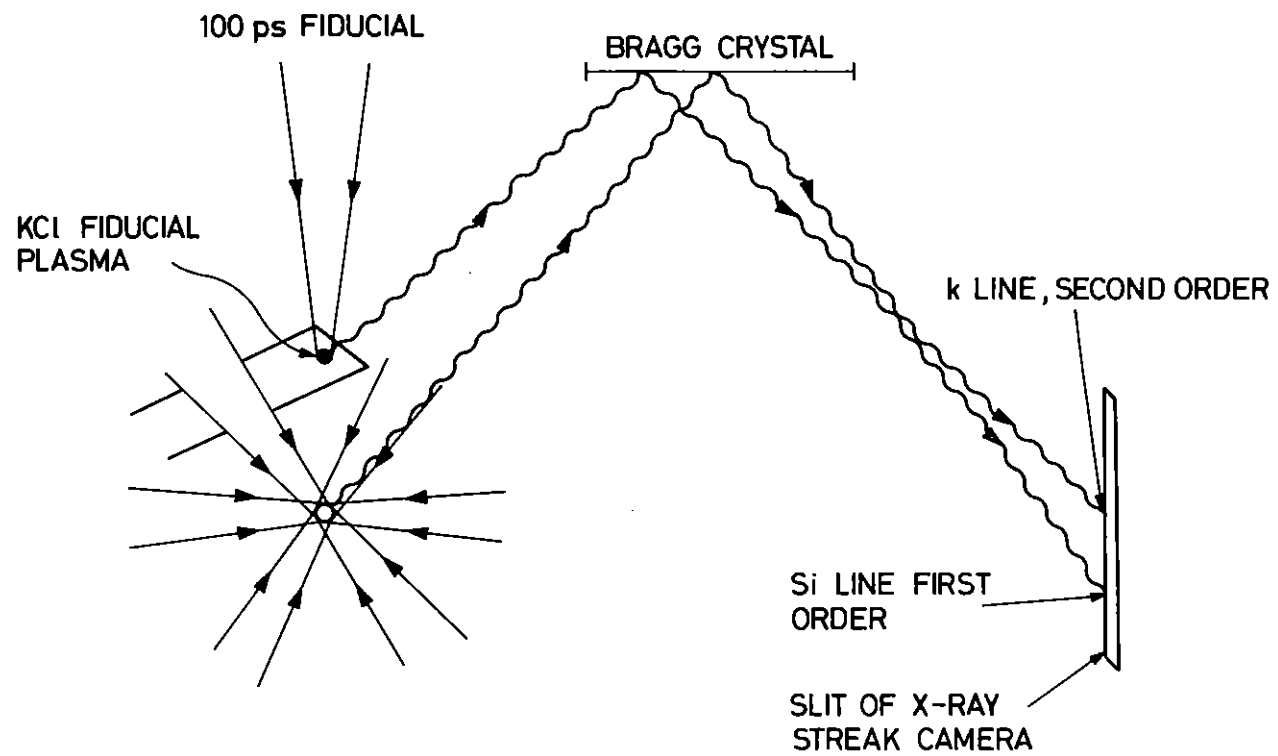


Fig. A1.15 Configuration for time resolved X-ray spectroscopy

(e) Time integrated crystal X-ray spectrograph.

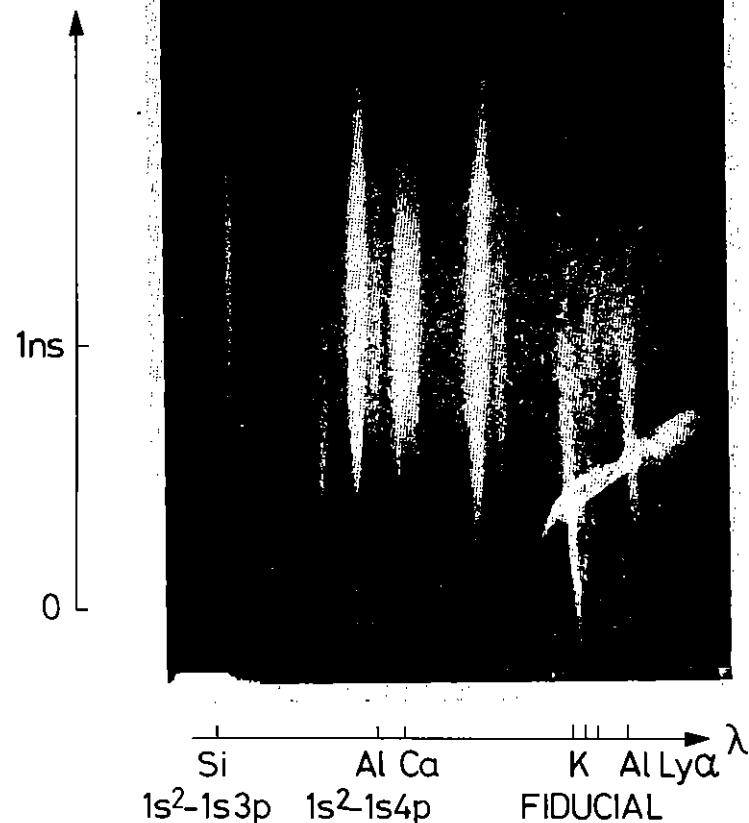
(f) A seven channel (filtered) X-ray diode array was used to measure the X-ray continuum. This diagnostic and inner shell line measurements were used to evaluate the role of hot electrons.

On most of the shots in this series, we attempted to have uniform illumination on the surface of the spherical targets. We were also concerned, however, with the role of refraction for rays with significant departure from normal incidence. In the early part of the experiment, we thus tried a variety of focal conditions as illustrated. The goal was to have as uniform illumination as possible while using relatively little overfocusing.

Uniformity of illumination was diagnosed with X-ray pinhole imaging, Faraday rotation probing and the previously mentioned interferometric probing. The pinhole imaging was used as a criterion of uniformity of energy deposition. The Faraday probe was used to measure the magnetic fields that result from nonuniformity in illumination. It was found that for $D/R \sim 2.25$ (tangential illumination) good uniform X-ray photographs (for a variety of X-ray energy channels from 1 to 10 keV) were obtained with negligible (large scale) magnetic field generation (< 100 kG). For uniform illumination, the absorbed fraction measured in these experiments (with ion calorimetry) was about 27%.

For incident irradiances of about 6×10^{14} W/cm², the hot electron temperature as measured from the X-ray continuum slope was about 2.5 keV. Inner shell (K_{α}) emission was measured from the calcium in the glass. In addition, K_{α} was observed from solid vanadium targets coated with various layers of plastic. From these measurements, we conclude that no more than 1% of the absorbed energy was present in hot electrons.

An example of a time resolved line spectrum is given in Figure Al.16. Clear delays can be seen between the time fiducial and the onset of aluminium



COATED SPHERE WITH KCl FIDUCIAL
(12260883)

emission and between the aluminium and silicon emission. Both theory and experiment show a very rapid rise (< 100 ps) for the onset of silicon and aluminium emission. The delays are thus well defined quantities. Average mass ablation rates can be obtained by taking these intervals as the time needed to ablate either the outer plastic or the aluminium and inner plastic layers.

In Figure A1.17 we plot the mass ablation rate versus absorbed irradiance. For any particular point the absorbed irradiance is taken as the average value during the delay (measurement) interval. The points on Figure A1.17 also represent a range of values of both irradiance and plastic thicknesses. The measurement intervals are thus distributed across various parts of the laser pulse and do not emphasise any particular part of the temporal profile.

Time resolved measurements of calcium emission from the glass were also made. For the normal tangential focused shots, a very small delay (< 100 ps) was observed between the turn on of silicon and calcium line emission. This would seem to indicate a relatively steep profile of the heat front.

Also shown in Figure A1.17 is the range of data in a recent study of mass ablation at KMSF (A1.27) and points taken at the University of Rochester (A1.28). The KMSF data agrees reasonably well with the present work while the Rochester points are somewhat higher. Reconciling this apparent discrepancy could be very important for laser fusion. At this point, it is not clear whether the higher fluxes would be a benefit or detriment. Enhanced penetration might lead to a form of preheat degrading the efficiency of laser driven implosions. In addition, the larger heat transport in the Rochester work seems difficult to explain by present theoretical methods.

Some of the differences between the Rochester work and the present study are:

- (a) $f/1$ focusing was employed in the present work while $f/3$ was used at Rochester.

- (b) Higher absorption was observed in the Rochester experiments.

- (c) Rochester work employed time integrated X-ray spectroscopy while the present work was time resolved.

As a point of comparison with the Rochester work, time integrated X-ray spectra were also measured in the present experiments. The assessment of transport using time integrated spectra is made by observing the reduction of intensity of the radiation of the aluminium or silicon as a function of the plastic thickness. In Figure A1.18, we plot some time integrated spectral results. We also show results from the Rochester experiments at a slightly different incident intensity. Since time delays were the primary diagnostic in the present work, a smaller range of plastic thicknesses were employed than in the Rochester experiment. In the present work however, there is clearly little difference in the fall of intensity between high and low Z radiators. A substantial difference was observed in the Rochester experiments.

A further indication of the differences between the Rochester experiments and the present work may be obtained from the small number of shots taken in the present work at non-tangential illumination. These shots yielded an apparently higher mass ablation rate, as shown in Table A1.1, although they were very close in irradiance and measured absorbed energy.

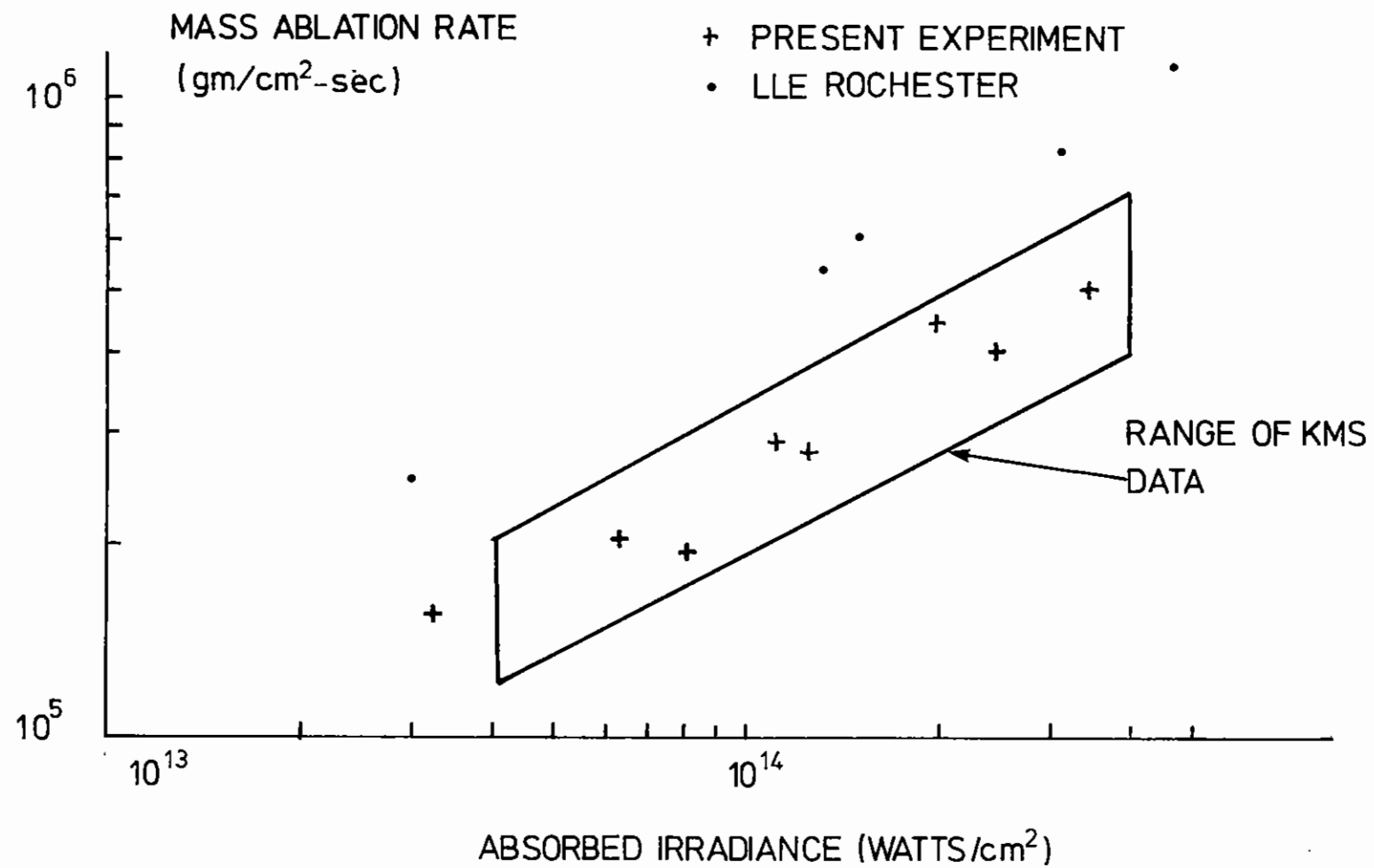


Fig. A1.17 Mass ablation rate as a function of irradiance during measurement interval.

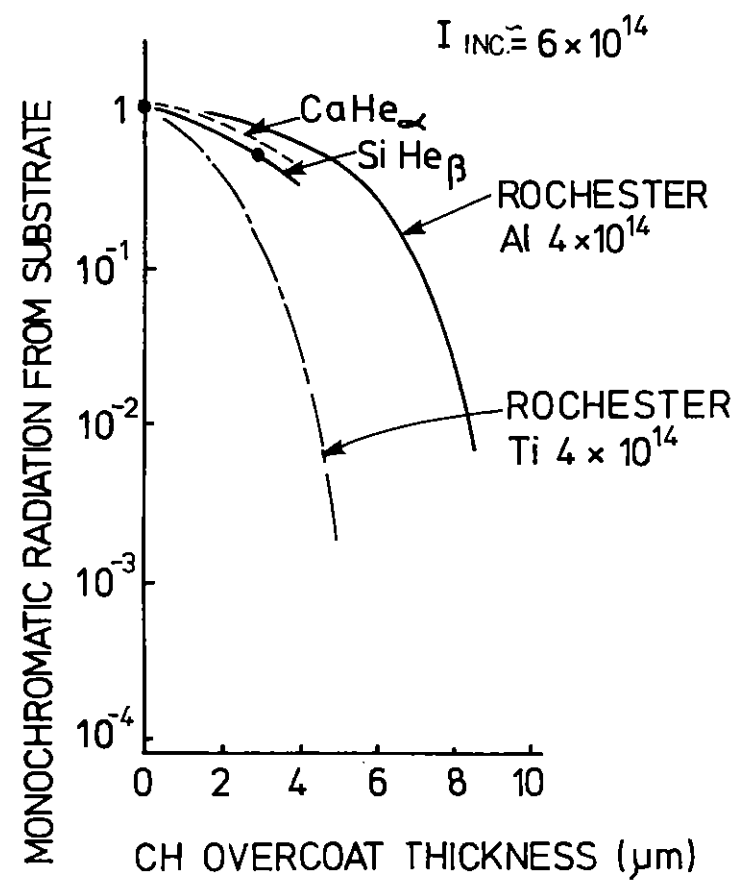


Fig. A1.18 Radiation intensity as a function of transport layer thickness

Table Al.1 Mass ablation for non-tangential illumination

	$\dot{m}$ gm/cm ² -s	Absorption	Incident Irradiance $\times 10^{14}(\text{W cm}^{-2})$
Tangential D/R $\approx$ 2.25	1.8	24%	2.8
Non-tangential 1.5 > D/R > 1.0	3.3	25%	3.3

The present experiments have been modelled with the Lagrangian hydrodynamic code LASNEX. With this modelling, thermal transport is treated as flux limited diffusion.

With the modest hot electron temperature, the non-local deposition of the hots is not dramatically different from that which would occur in a single group transport model.

In Figure Al.19, we show the result of a LASNEX calculation of emission. The delay between aluminium and silicon emission at $3.5 \times 10^{14} \text{ W cm}^{-2}$ absorbed irradiance should be compared with 440 ps observed in the experiments. In making such direct comparisons we find that the best results are obtained when a flux limit of about 0.08 is used. This flux limit also gives the best comparison with measured expansion profiles as shown in Figure Al.20.

In conclusion, we find that the thermal electron transport in these spherical target experiments is best described by near classical modelling. The observation of little difference between emission from high and low Z elements is also consistent with this picture. Some clue as to the difference between the present work and the higher flux observed in the Rochester experiments may be given by the tighter focus shots described in Table Al.1. Further analysis of this discrepancy is needed.

Al.6 Heat Flow and Ion Acoustic Turbulence in a Laser-heated Z-pinch Plasma

A E Dangor, A K L Dymoke-Bradshaw, M Lewis (IC), D A Pepler and W T Toner (RAL)

An experiment was carried out to investigate the heat flow and ion acoustic turbulence generated when a large temperature gradient is induced in a plasma. Such a gradient was created by focusing the $\lambda = 1.05 \mu\text{m}$ laser beam into a fully ionised hydrogen Z-pinch plasma. The plasma parameters were then measured by collecting and analysing 90° scattered light from a $\lambda = 0.53 \mu\text{m}$ probe beam (see Figure Al.22).

A trial experiment had been carried out to look at the heat flow (see Section 5.9 of the RAL 1983 Annual Report to the LFC), but this was limited in that it was not possible to determine the relative positions of the heated and scattering volumes with sufficient accuracy. This was rectified by the inclusion of an alignment system which had a reproducibility of $< 10 \mu\text{m}$. The collection optics for the scattered radiation were also improved.

Figure Al.23 shows the experimental layout. The $\lambda = 1.05 \mu\text{m}$ heating beam used was a 2 ns flat-topped pulse, focused to a spot of $\sim 100 \mu\text{m}$ diameter on the axis of the pinch. The $\lambda = 0.53 \mu\text{m}$ probe beam was also of 2 ns duration. The light scattered through 90° was collected, and the output spectrum is of the form shown in Figure Al.24, with the exact shape being dependent on the electron temperature and density. There are two parts to the spectrum: broad wings due to the electrons, and a much stronger, narrow ion feature centred on the probe beam wavelength. The scattered beam was split into two, with $\sim 10\%$ of the total light going to a high dispersion spectrometer, and the rest to a second spectrometer with lower dispersion. The output from the first spectrometer had a spectral range of $\sim 10 \text{ \AA}$ and was centred on the probe wavelength, thus giving a high resolution ($\sim 1 \text{ \AA}$) image of the ion feature. The whole spectrum was recorded from the second spectrometer, with a spectral range of $\sim 250 \text{ \AA}$, and $\sim 4 \text{ \AA}$ resolution. Both

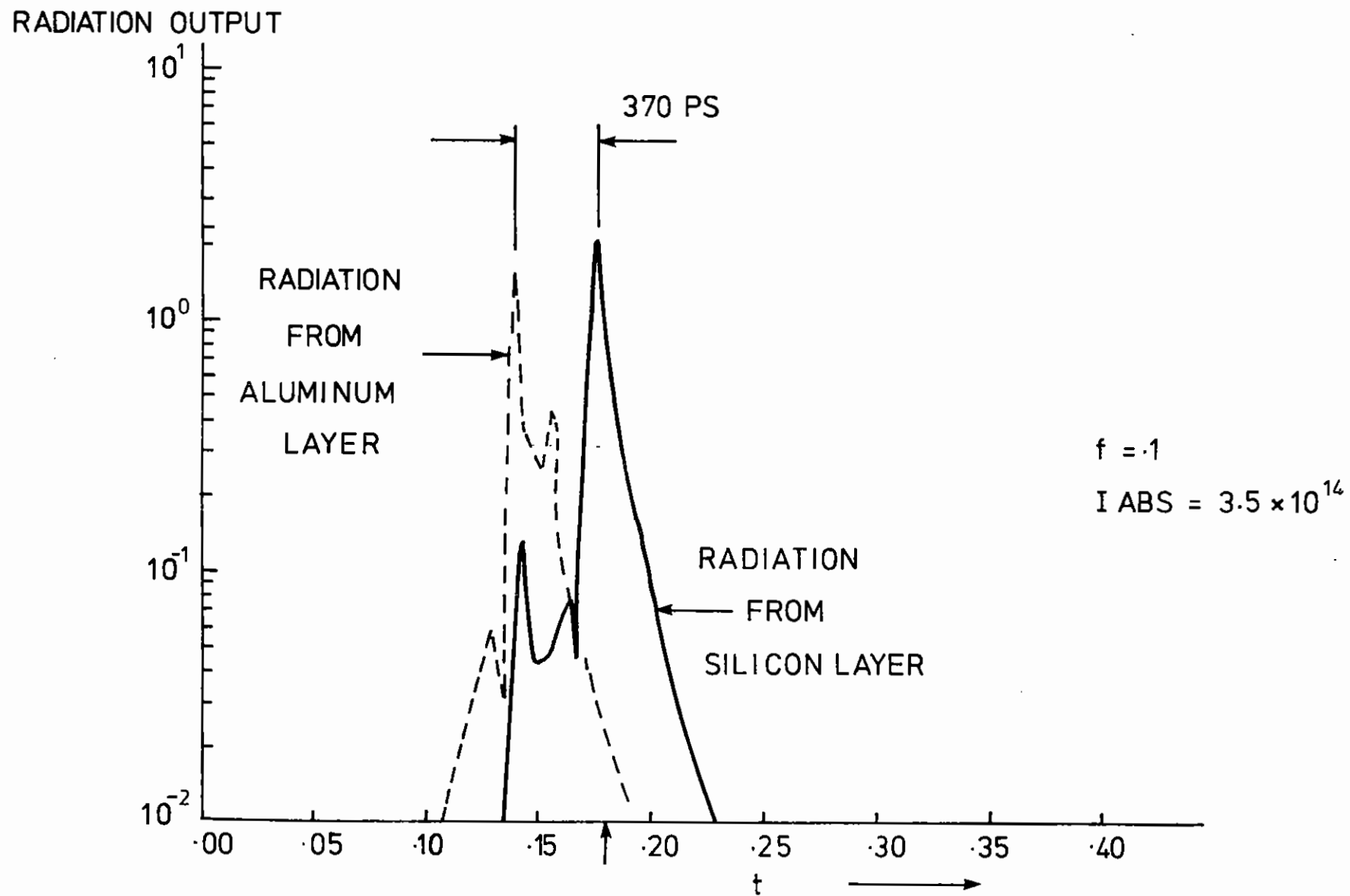


Fig. Al.19 Comparison between lasnex hydro simulation and experiment.

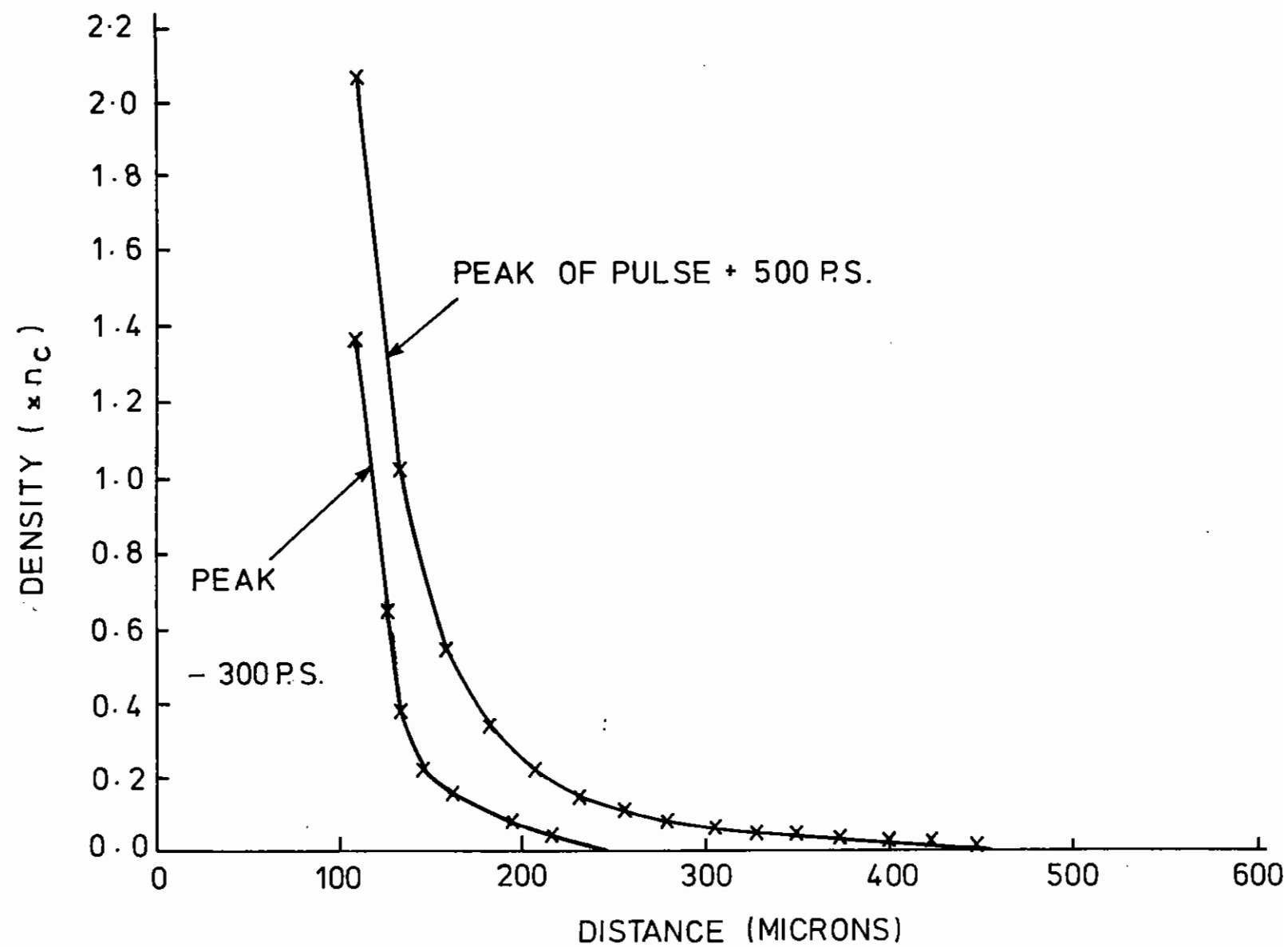


Fig. A1.20 Density profiles from interferometric probing.

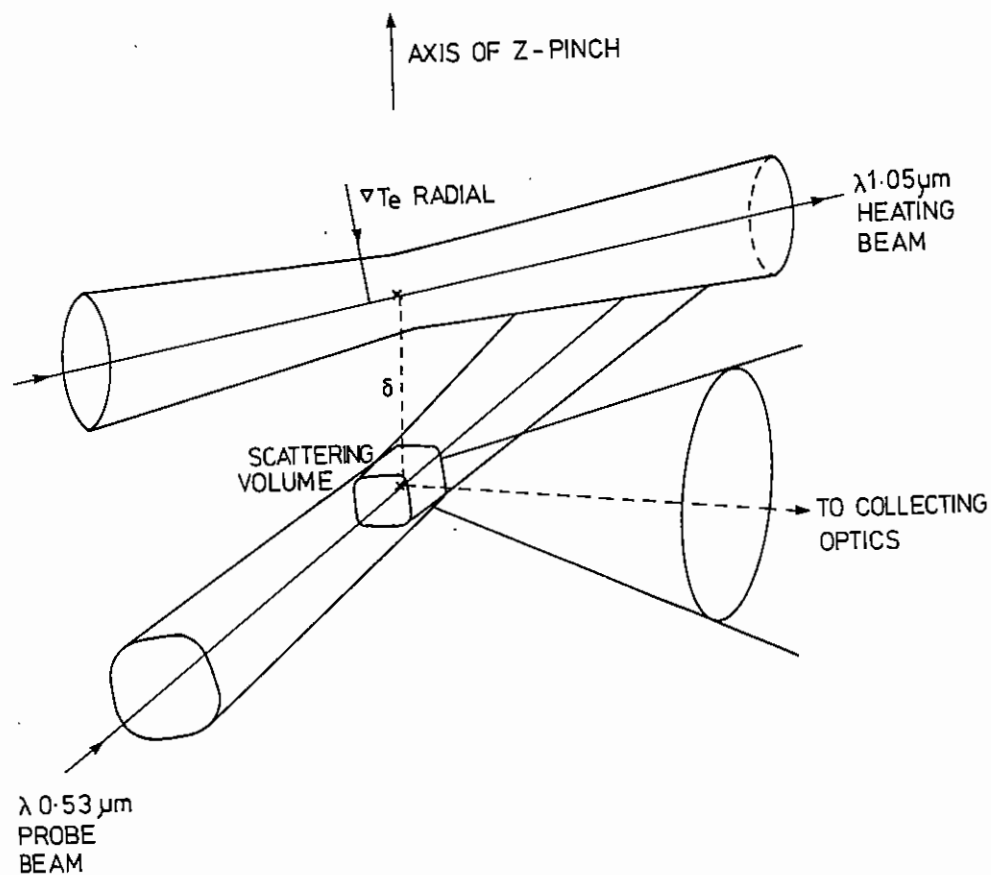


Fig. A1.21 Spatial relationship of laser beams: probe beam at 45° to heating beam, with spatial displacement of δ ; optics set to collect radiation scattered through 90° .

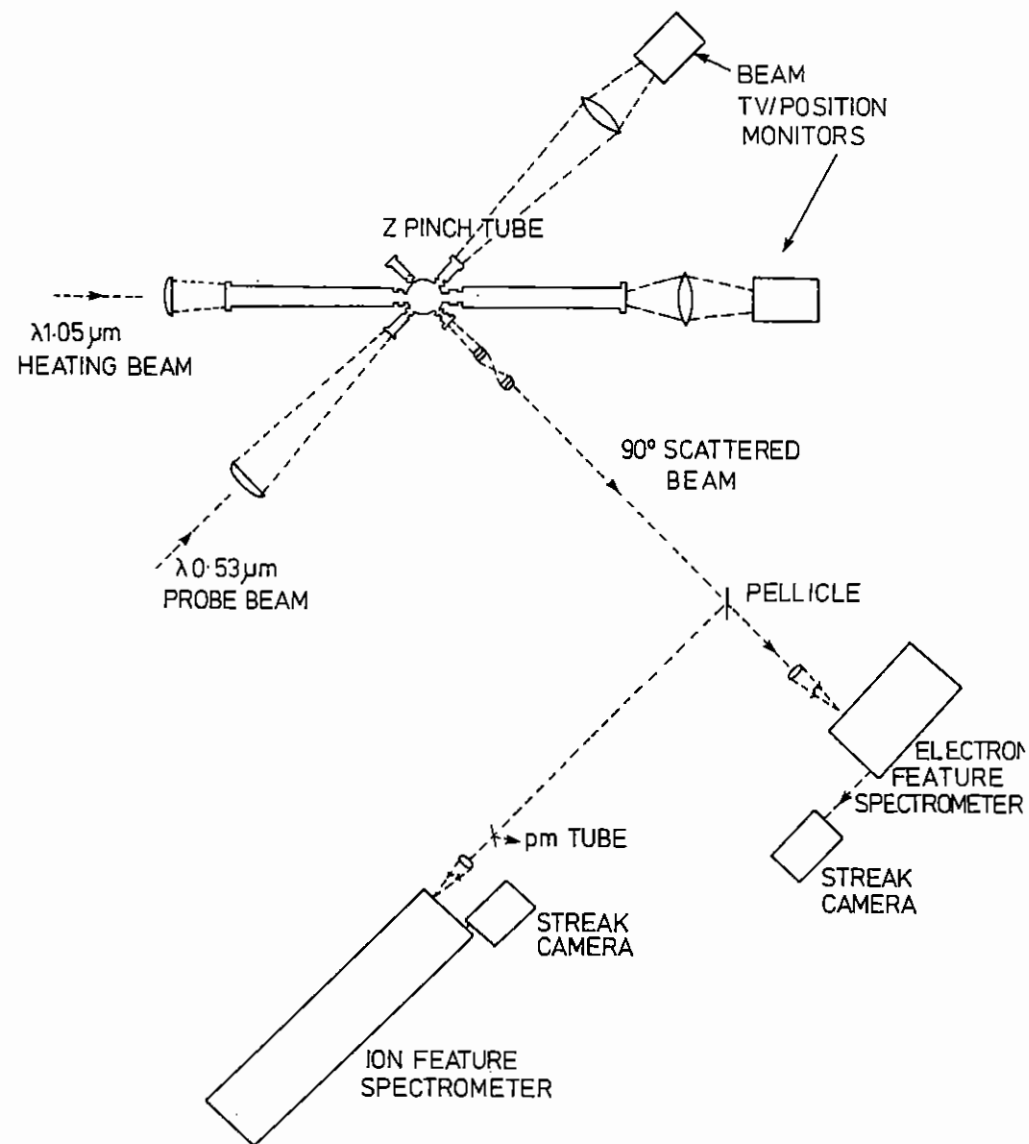


Fig. A1.22 Experimental layout.

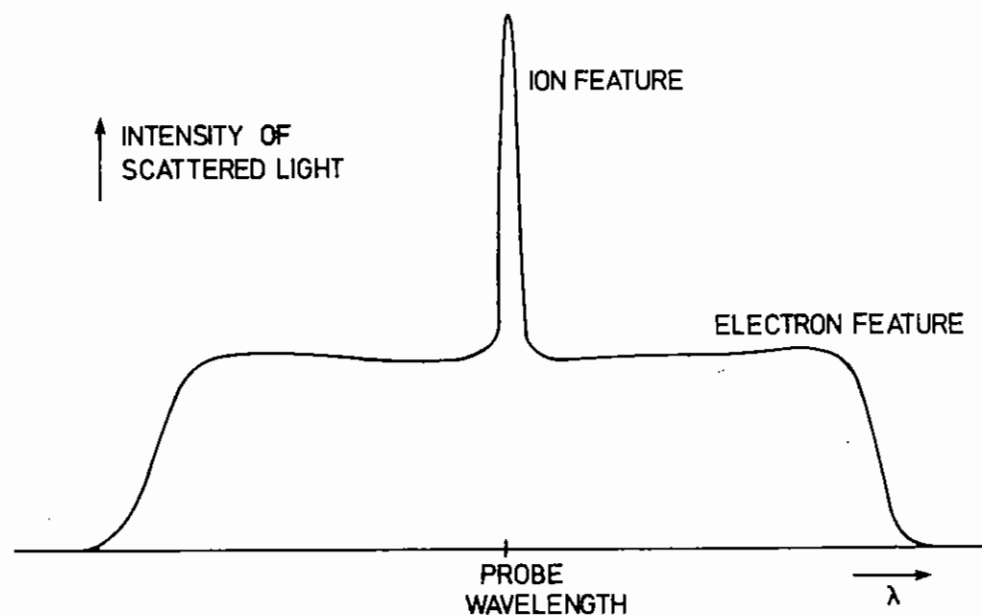


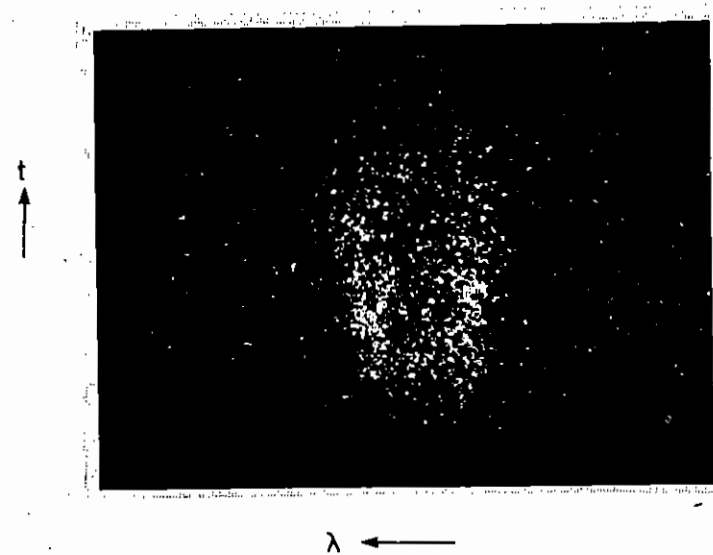
Fig. A1.23 Generalised schematic of the scattered spectrum, with an intense ion feature (width $\sim$ few $\text{\AA}$) and a broad ($\sim$ few hundred $\text{\AA}$), low intensity electron feature.

of these spectra were recorded on film using streak cameras, with a time resolution of ~ 0.4 ns.

It was possible to change a number of parameters throughout the experiment, thus allowing a full study to be made of the plasma conditions. For example (a) the position of the heating beam could be moved, with respect to the diagnosed volume, by up to several hundred microns both vertically and horizontally, thus allowing the relative directions of k_s and ΔT to change; (b) the relative timing of the heating and probe beams could be varied, and (c) it was possible to trigger the Z-pinch at any time up to 6 μs before the laser was fired, thus enabling the initial n_e and T_e to be varied.

Ion and electron spectra were recorded for a number of these conditions, and typical outputs are shown in Figure A1.25. As the ion feature is much brighter per unit bandwidth than the electron feature, a strip of neutral density filter was placed in the focal plane of the spectrometer centred on the ion feature to prevent the film saturating. This shows up in the photograph as a dark strip, with the ion feature showing through in the centre. Figure A1.26 is a series of microdensitometer traces of these features as a function of time. The ion feature clearly shows that ion acoustic turbulence is present, and analysis of the electron feature gives $n_e \sim 10^{17}$ and $T_e \sim 25$ eV. Many shots exhibit an asymmetry in the electron feature, and some also in the ion feature. In the latter, this is probably due to a relative electron-to-ion drift velocity, which produces ion acoustic waves at ω_{pi} if the Landau damping is low (ie T_e/T_i is large). Asymmetries in electron features are harder to explain, but may be driven by the fast streaming and slow counter-streaming electrons. The latter form a space charge neutralising return current. Analysis of all the data is continuing at present.

(a) ION FEATURE



(b) ELECTRON FEATURE

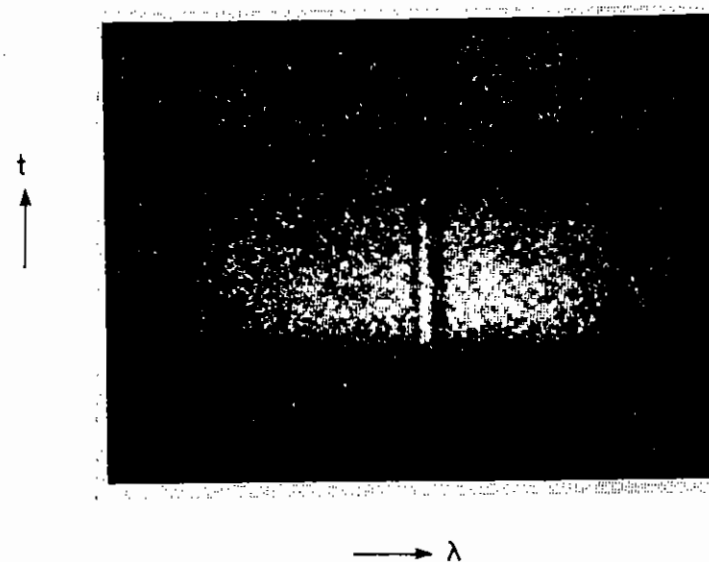


Fig. A1.24 Streak photographs of the ion and electron features, with the heating and probe beams coincident in time and space.

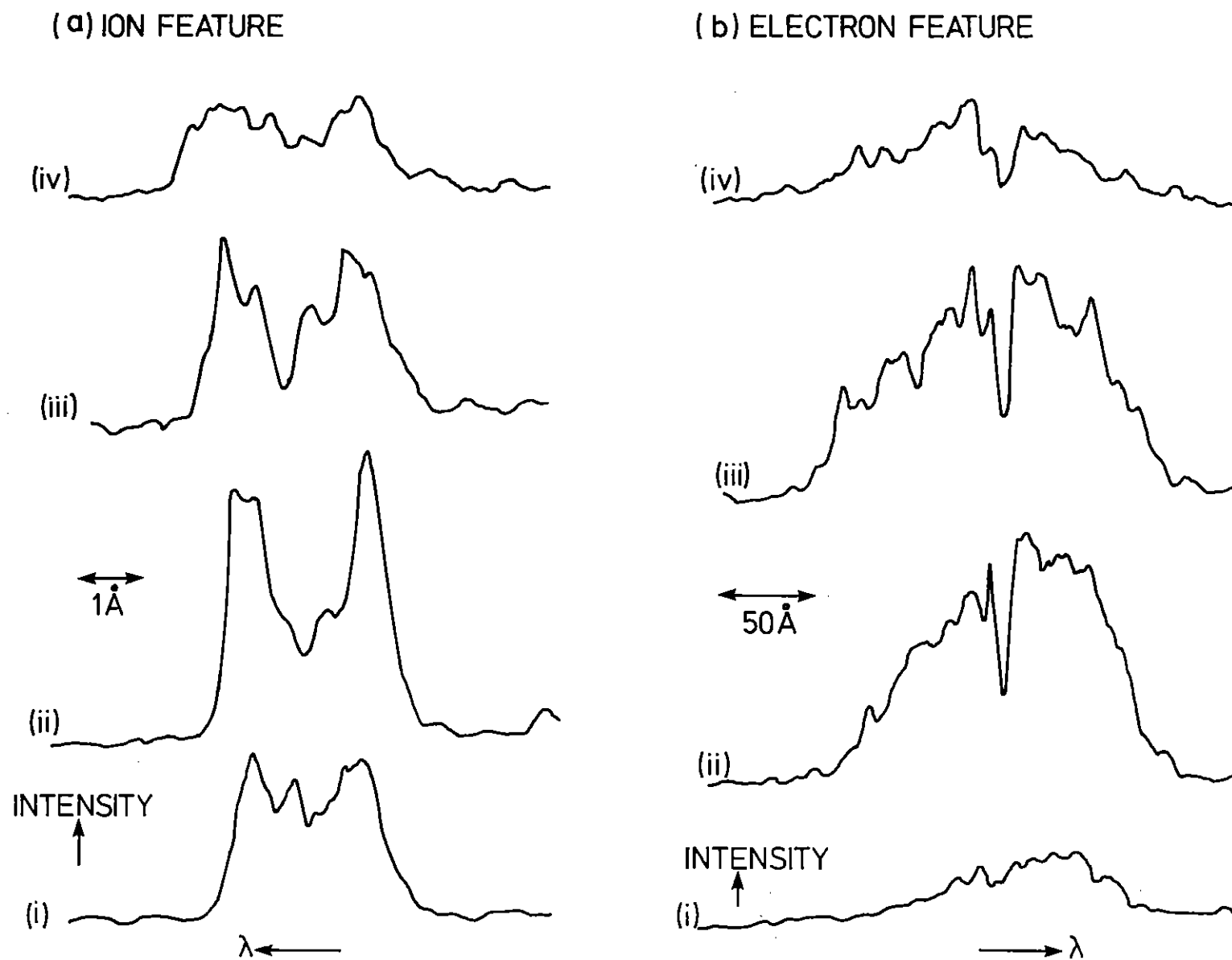


Fig. A1.25 Microdensitometer traces of the ion and electron features, with the heating and probe beams coincident in time and space. Traces (i) to (iv) show the time variation for the duration of the laser pulse.

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A2 LASER PLASMA COMPRESSION AND DYNAMICS

A2.1 Introduction

T A Hall (Essex) and A J Cole (RAL)

The results reported in this section cover all the implosion and ablative acceleration experiments whose main aim is to improve our understanding of the processes occurring in laser driven compression of microballoon targets.

The highlight of the years work was the systematic study of ablative implosions with green laser light (section A2.2). This "campaign" experiment, comprising of a logical progression from uniformity studies (section A2.2.3) through to X-ray radiographic and emission diagnosis measurements (sections A2.2.4, 5 and 6), was carried out by experts in each field from the various university groups active at the C.L.F., together with RAL staff. Careful planning of the experiments, including extensive simulations of target performances (Section A2.2.2) was an important contribution to the success of the campaign.

Several new techniques for observing implosion experiments were tried out during the reporting period including point-source radiography (section A2.3) and the x-ray framing camera developed at Essex university (Section A2.4).

Work continued on the Rayleigh-Taylor instability (Section A2.5) and also on thermal smoothing effects in targets irradiated with third harmonic (0.35 μ m) laser light (Section A.2.6).

A2.2 Systematic Study of Six Beam 0.53 μ m Ablative Implosions

A2.2.1 Introduction

M H Key (RAL), J D Kilkenny (IC) and C L S Lewis (QUB)

A long term theme of the work of the scientific group for 'laser plasma compression and dynamics' has been the endeavour to understand and optimise ablatively driven implosions.

The major variables are the target parameters viz radius r , wall thickness Δr and wall material (density ρ and atomic number Z) and the irradiation parameters viz laser wavelength λ , irradiance I and pulse duration τ .

Fairly good understanding of some of the limiting factors has been obtained from experiments and calculations isolating specific processes. Pressure multiplication is proportional to $r/\Delta r$ but study of R-T instability has shown growth rates of about 50% of the classical value suggesting a growth exponent $\gamma_t \sim (r/\Delta r)^{1/2}$ in an implosion and therefore an upper limit $r/\Delta r \dagger 10$ (A2.1).

If the power P of the laser is limited then the irradiance is linked to the target size ($P \sim 4\pi r^2 I$). Consideration of hot electron generation suggests that there is an upper limit to $I\lambda^2$ arising from the range of hot electron preheating $(\rho\lambda)_h$ which is proportional to $I\lambda^2$. Targets with $r/\Delta r \approx 10$ and of a size matched to the CLF Nd glass laser have $\rho \Delta r \sim 10^{-3} \text{ g cm}^{-2}$. For hot electron preheating to be less than shock preheating there is a limit $I\lambda^2 \dagger 10^{14} \text{ W cm}^{-2} \mu\text{m}^2$.

Preheating by soft X-rays is also important and may cause a loss of pressure and density in implosions which can only be overcome by increasing the wall thickness of the target and thus, for constant $r/\Delta r$, the radius. Since the laser energy E_L is proportional to r^3 this is an expensive requirement. Conditions where soft X-ray preheating does not give a significant temperature rise relative to shock preheating in the inner half of the pusher are obtained for $\rho\Delta r > 10^{-3} \text{ g cm}^{-2}$ for low Z shells (A2.2, A2.3).

The laser wavelength λ is a key variable. At $I\lambda^2 \sim 10^{14} \text{ W cm}^{-2} \mu\text{m}^2$ and for absorption fraction $A = 0.3$ the ablation pressure $P_a/1 \text{ Mbar} = 4(\lambda/1 \mu\text{m})^{-2}$ (A2.4). Thus shorter wavelengths are advantageous.

Problems at shorter wavelength include reduced thermal conduction smoothing of spatial pressure variations arising from spatial irradiance variations of scale length L . For the relevant irradiance range an approximate thermal smoothing factor f for an implosion is given by $f \sim 0.1(L/r)(r/\Delta r)^{1/2}(\lambda/1 \mu\text{m})^{-0.9}$ (A2.5). In multibeam irradiation schemes individual focused beams have typically 2:1 intensity variation (A2.6). The need for pressure uniformity of a few percent and thus $f < 0.1$ may mitigate against wavelength $\lambda < 0.5 \mu\text{m}$. There is considerable uncertainty on this point however due to lack of good experimental data.

The above discussion also suggests that there is negligible thermal smoothing on the spatial scale $L/r \sim 1$ associated with the pattern of beam overlap in multibeam irradiation. There is therefore a requirement for a few percent uniformity of the irradiation itself. Calculations indicate that six beam symmetric irradiation is capable of approximately 15% uniformity, with 12 beams giving an improvement to about 3% (A2.7).

It was decided therefore to conduct a careful study of the performance of optimised targets driven by six beams in preparation for future experiments at the CLF using 12 beams.

The optimised conditions would be $r/\Delta r = 10$, $\lambda = 0.5 \mu\text{m}$, $I \sim 4 \times 10^{14} \text{ W cm}^{-2}$ which with the 300 J in one nsec available from the six beam VULCAN laser set the target radius at approximately 80 μm .

A coordinated series of experiments was planned to investigate the performance of optimal six beam ($f/1$) ablative compression in which the initial phase was designed to optimise the uniformity of irradiation and measure the absorption fraction (section A2.2.3). Images of harmonic emission at 2ω and $3/2\omega$ and of X-ray emission were used to determine the conditions of most uniform irradiation varying the overlap of the beams on

the target (d/r where d is the distance of the focal point beyond the target centre).

It was concluded that best uniformity was obtained for $d/r \sim 4.5$ as suggested also by calculation (A2.7). The absorption was $A \sim 0.2$. The X-ray data also gave a quantitative estimate of the ratio of maximum to minimum ablation rate of 1.25 ± 0.05 for the three beam overlap points relative to the 1 beam centres.

In the second phase of the experiment (A2.2.4) pulsed X-radiography was used to diagnose the symmetry of implosions and the maximum density.

The choice of targets for this work was influenced also by the requirement of the diagnostic method. It is advantageous to use a low Z ablator giving low intensity of X-rays from the ablation plasma relative to the backlighting source. Maximum stability against R-T growth is obtained with a low density low aspect ratio homogeneous shell. Maximum X-ray probe penetration is obtained using a low Z shell. The choice was therefore a homogeneous polymer shell. The gas fill density should ideally be large enough that the compression ratio in a perfectly symmetric implosion would not exceed the practical limit of about 1000 $\times$ expected to arise from imperfect sphericity. Freon gas fill at densities $\rho \sim 10^{-2} \text{ g cm}^{-3}$ was therefore envisaged (in the event empty polymer shells were used in the experiment for reasons of target availability).

Practical consideration of backlighting source intensity relative to target self emission also set a limit to the irradiance on target. Typical conditions chosen were $I = 2 \times 10^{14} \text{ W cm}^{-2}$, $r = 80 \mu\text{m}$, $r/\Delta r = 10$ with an implosion time and pulse duration of about 1 ns. Simulation of the implosions indicated real compressed density of about 15 g cm^{-3} with the effects of radiative preheat being relatively unimportant (A2.3).

Observation of the uniformity of the initial phase of implosion verified the conclusion that for $d/r \sim 4.5$ the pressure maxima were at the 3 beam overlap and minima at the individual beam centres with a pressure ratio of about 1.3:1. Average density in the core at peak compression was measured to be

10 gcm⁻³. This is close to the 1D modelling result for conditions of fairly negligible preheat induced decompression of the imploding shell and suggests that the optimisation discussed earlier is soundly based. Our earlier experiments with high aspect ratio shells and non optimised conditions all showed much lower density and strong departure from this 1D ideal performance (A2.6).

In the third phase of the experiment the objective was to produce 'hot' ablative implosions in which the core temperature was sufficient for X-ray and thermonuclear emission diagnostics (A2.8).

Target designs are discussed in A2.2.2 and these were of two types.

For X-ray emission the target design was based on a homogeneous plastic shell of low X-ray opacity with a Freon fill gas as discussed above. The Freon was chosen to give line radiation from Cl XVI of sufficient photon energy to escape through the low Z pusher at peak compression ($\rho\Delta r \sim 10^{-2}$ gcm⁻²). It is difficult however to satisfy all the constraints for optimal implosion and obtain the necessary core temperature. In practice reduction of the aspect ratio constraint was needed to achieve a higher implosion velocity and core temperature. The effect was to reduce density through preheat induced decompression and/or RT instability. Results are discussed in A2.2.5.

For thermonuclear emission the chosen targets were D-T filled at 10 bar pressure. 10 bar Ne was added to raise the density to 10⁻¹ gcm⁻³ and reduce the compression ratio. A glass shell was needed to contain the D-T and a homogeneous glass shell was preferred for better stability against R-T growth. The higher shell density implies a slower implosion velocity for a given aspect ratio. It was therefore necessary to reduce the shell thickness and increase the aspect ratio above 10 in order to achieve sufficient implosion velocity and temperature for fusion.

Typical conditions were $I = 5 \times 10^{14}$, $\tau = 0.5$ ns, $A = 0.2$, $r = 80$ μ m, $\Delta r = 1$ to 3 μ m. The experiments are described in section A2.2.6. Neutron yield and temperature decreased with wall thickness increase ranging from 10⁶, 400

eV to 7×10^3 , 260 eV (Figure A2.45). The temperature is proportional to the inferred value of momentum flux density ρv^2 in the imploding shell deduced from the shell parameter and ablation pressure, as expected for an ablative implosion.

Measurement of the value of $\rho\Delta r$ at the time of thermonuclear burn was made from the energy loss of 3.5 MeV DT α particles and 2 MeV DD protons which have ranges of approximately 2×10^{-3} gcm⁻² and 2×10^{-2} gcm⁻² respectively. Energy measurements were made by analysis of etched tracks in CR39 and even though DD proton yields are 10⁻² of DT α yields, good data were obtained for both.

The $\rho\Delta r$ at the time of burn was typically 10 x the initial $\rho\Delta r$ of the shell with maximum $\rho\Delta r \sim 10^{-2}$ gcm⁻² obtained for the largest (3 μ m) wall thickness. This is consistent with density of 10 gcm⁻³ and the corresponding ion temperature is 200 eV.

The CR39 technique was seen to be an extremely valuable and accurate diagnostic with interesting possibilities for deducing the source volume and uniformity of the pusher from the energy spread of the α and p emission as discussed in section A2.2.6.

Future work will give a direct comparison between 6 beam and 12 beam target irradiation.

A2.2.2 Target Design Calculations

B J MacGowan (IC) and S J Rose (RAL)

Prior to the start of the Campaign Experiment attention was focussed on the design of ablatively - driven, single - shell microballoons for the x-ray emission (A2.2.5) and thermonuclear spectroscopy (A2.2.6) sections. Glass microballoons filled with a mixture of Neon and DT were chosen for the thermonuclear spectroscopy experiments whereas Freon-filled plastic shells were used in the x-ray emission work. In each case so as to produce a sufficiently high output, it was intended that the targets produce a high

maximum gas temperature of at least several hundred eV. Several other constraints were taken into consideration. Firstly the aspect ratio must be low enough (less than about 10) that Rayleigh-Taylor instability does not break up the shell during the implosion. Also the target wall must be sufficiently thick that fast-electron or radiative preheat does not greatly affect the implosion. 1D MEDUSA simulations were used to design the targets, shown in Table A2.1, which fit as far as possible these requirements. They were designed assuming the output of green laser light to be 400 J which together with 77% transmission through the lenses and an assumed absorption of 35% results in approximately 100 J absorbed.

Fast electron and radiation transport were each included using a multi-group approximation. Care must be exercised, however, in interpreting the results of such calculations. For example, the code predicts an unrealistically high ion temperature when the spherically convergent shock reaches the centre. Also because targets could only be manufactured with low gas densities, the resulting maximum compressions predicted by 1D simulations (over 3000) are unlikely to be achieved in practice. Although it is possible to circumvent these difficulties by ignoring those parts of the simulation which are thought to be unrealistic, there still remains considerable uncertainty in the predicted target performance.

Table A2.1 Calculated target performance

Irradiance (Wcm ⁻²)	Pulse Length (ps)	Shell Outer Radius (μm)	Shell Thickness (μm)	Shell Material	Initial Gas Fill	Maximum Gas Temp.(ev)	Neutron Yield
9 x 10 ¹⁴	500	73	10	Plastic	2 bar Freon (10 ⁻² gcm ⁻³)	400	-
1.6x10 ¹⁵	500	61	5	Glass	10bar Neon + 10bar DT	560	10 ⁵

A2.2.3(a) X-ray Studies of Illumination Uniformity

J Wark, J D Kilkenny, B J MacGowan (IC)

(i) Introduction

The illumination uniformity was studied by recording the variation in brightness of X-rays emitted from irradiated microspheres. The X-ray images were recorded using two orthogonal pinhole cameras filtered with 25 μm Be +5μm Al. The cameras were positioned halfway between two of the six beams; thus on the horizon of the sphere were both double and triple beam overlaps, as well as two single beam positions (as shown in Figure A2.1).

In order to enhance the effects of any non-uniformities in the illumination the solid glass spheres of 70μm diameter were coated with 3-6μm of plastic. Thus in regions with higher absorbed energy the plastic would be ablated more quickly, revealing the brighter X-ray source of the glass. The Al

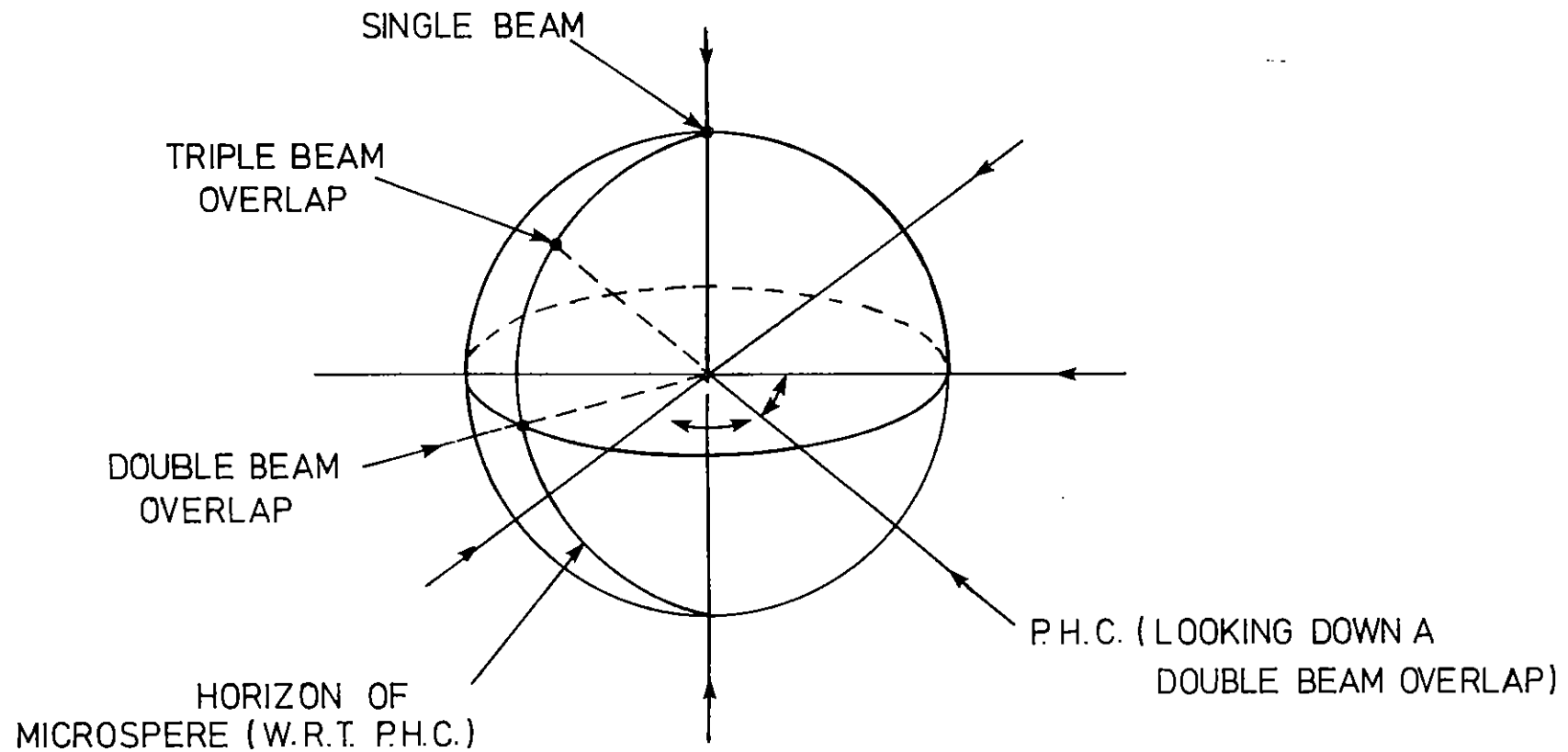


Fig. A2.1 Position of X-ray pinhole camera showing that on the horizon are triple and double beam overlaps, as well as single beam positions.

filtering was used so that the relatively long lived Silicon He α line was suppressed, and the X-rays from the glass closely followed the laser pulse. The fractional absorbed energy was also monitored using six ion calorimeters.

Over 20 shots were fired, the beams being focussed between 2.5 and 5.5 target radii beyond the microsphere to look for focussing dependence of the uniformity, at a typical incident irradiance of $5 \times 10^{14} \text{ W cm}^{-2}$. In order to obtain quantitative information about uniformity, it is necessary to know the relative brightness of glass and plastic X-rays with the filtering used. Figure A2.2 shows the average relative X-ray yield against plastic thickness for shots with similar absorbed irradiance. With $6 \mu\text{m}$ of plastic, burn through to the glass is not achieved, and it is not until the plastic thickness is reduced to $\sim 4 \mu\text{m}$ that the X-ray yield increases. From this we can estimate the mass ablation rate to be $\sim 7 \times 10^5 \text{ gm cm}^{-2} \text{ s}^{-1}$, which is about 30% higher than previous measurements using layered targets (A2.4). This is probably due to the lower X-ray energy used in this method, as the heat front has a finite gradient.

(ii) Qualitative Results

Typical results are shown in Figures A2.3 and A2.5 with their associated beam balances (Figures A2.4 and A2.6). It can be seen that with good beam balance ($\pm 2.5\%$), X-ray emission is dominant at the three beam overlap positions, with lowest yield at single beam points. However, it was found that if the beam balance was bad ($\pm 7.5\%$), no such structure could be seen, the maximum and minimum being apparently random. Only three shots were obtained with beam balance better than $\pm 3\%$, all at a defocus of 4.5 radii, and all of these shots show that three beam overlap positions have the highest X-ray yield. Little information can be extracted from the shots with poorer beam balance.

(iii) Quantitative Results

Assuming that the X-ray yield from the plastic and glass is proportional to the time the laser spends burning through the material, we can write:

$$I = A (0.6 - T) + BT$$

where I is the relative X-ray intensity, T the time the laser pulse spends burning through the glass in nanoseconds, and A and B are constants determining the relative brightness of plastic and glass X-rays respectively. From Figure A2.2 we see that at $T = 0$

$$I = 117 \pm 8 = 0.6A$$

Therefore $A = 195 \pm 13$ and at $T = 0.6$ (ie. no burn through to the glass)

$$I = 11 \pm 1 = 0.6B$$

Therefore $B = 18.3 \pm 2$

Using these values for the constants A and B it is possible to convert X-ray intensities to the time the laser beam spent in the glass, and hence determine relative differences in the mass ablation rate around the horizon of the microsphere.

Taking the average of the three results used, the mass ablation rate is ($25 \pm 5\%$) higher at a triple beam overlap than at a single beam position. This experimental result agrees qualitatively with the simple absorption model described by A R Bell in (A2.7) for an angular absorption dependence of $\cos \theta$ or $\cos^2 \theta$. The quantitative results cannot be directly compared as the effects of thermal smoothing in spherical geometry are not known, though planar smoothing results indicate that in the irradiance regime used some smoothing would take place (A2.9).

(iv) Uniformity of Illumination Control

Figure A2.7 shows the variation in absorbed energy with defocus as measured with the ion calorimeters. The energy falls off from $\sim 33\%$ at obscuration ($d/R = 2.25$ with $f/1$ lenses) to $\sim 17\%$ at $d/R = 5.0$. This is not merely a geometric effect (ie. less energy absorbed at higher d/R as less energy hits the target); the ratio of absorbed energy at $d/R = 2.25$ to that at $d/R = 5$ is ~ 2 , whereas by geometrical effects alone it would be ~ 5 . This is

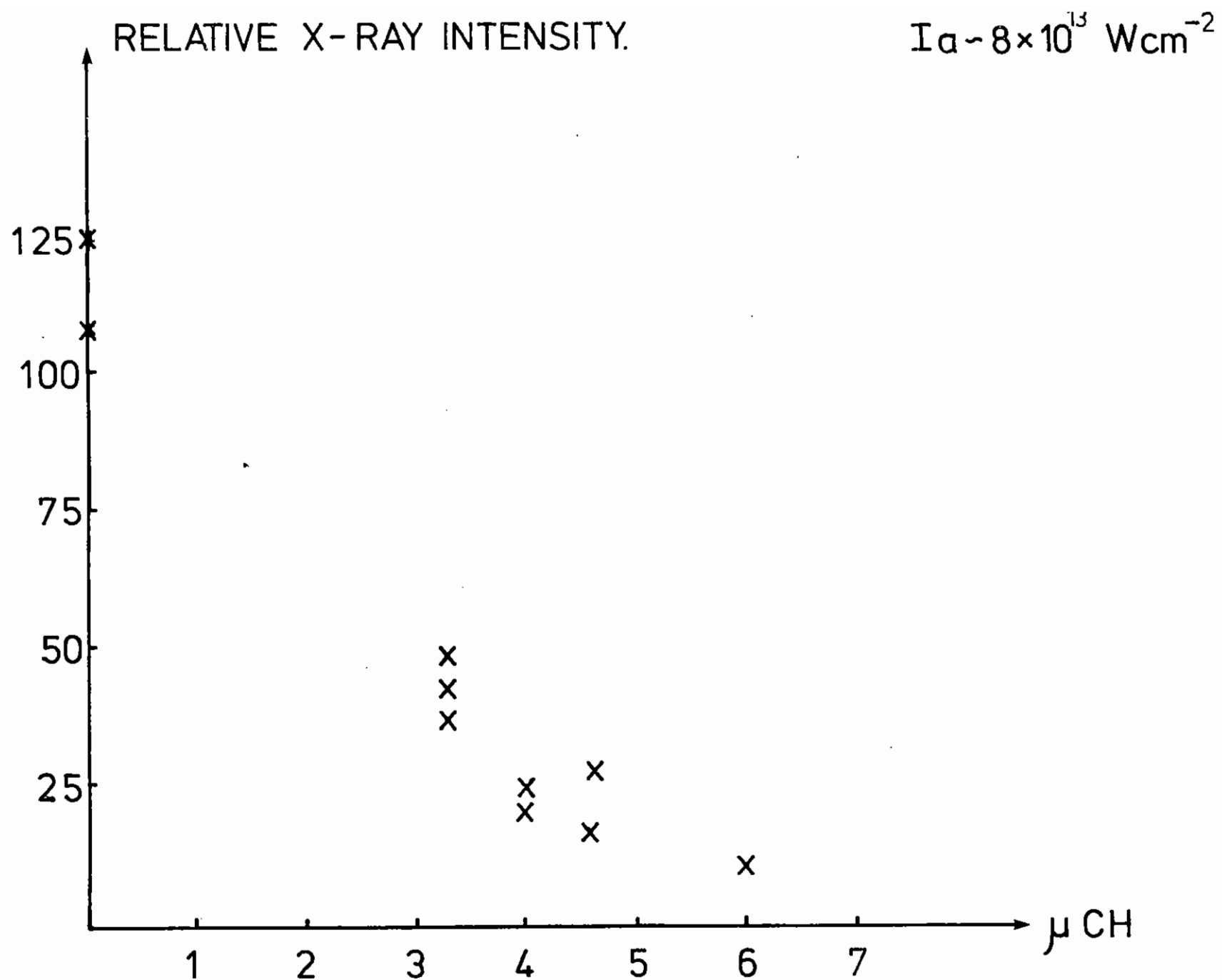


Fig. A2.2 X-ray intensity (arbitrary units) versus plastic coating thickness for an absorbed irradiance of $\sim 8 \times 10^{13} \text{ W cm}^{-2}$.

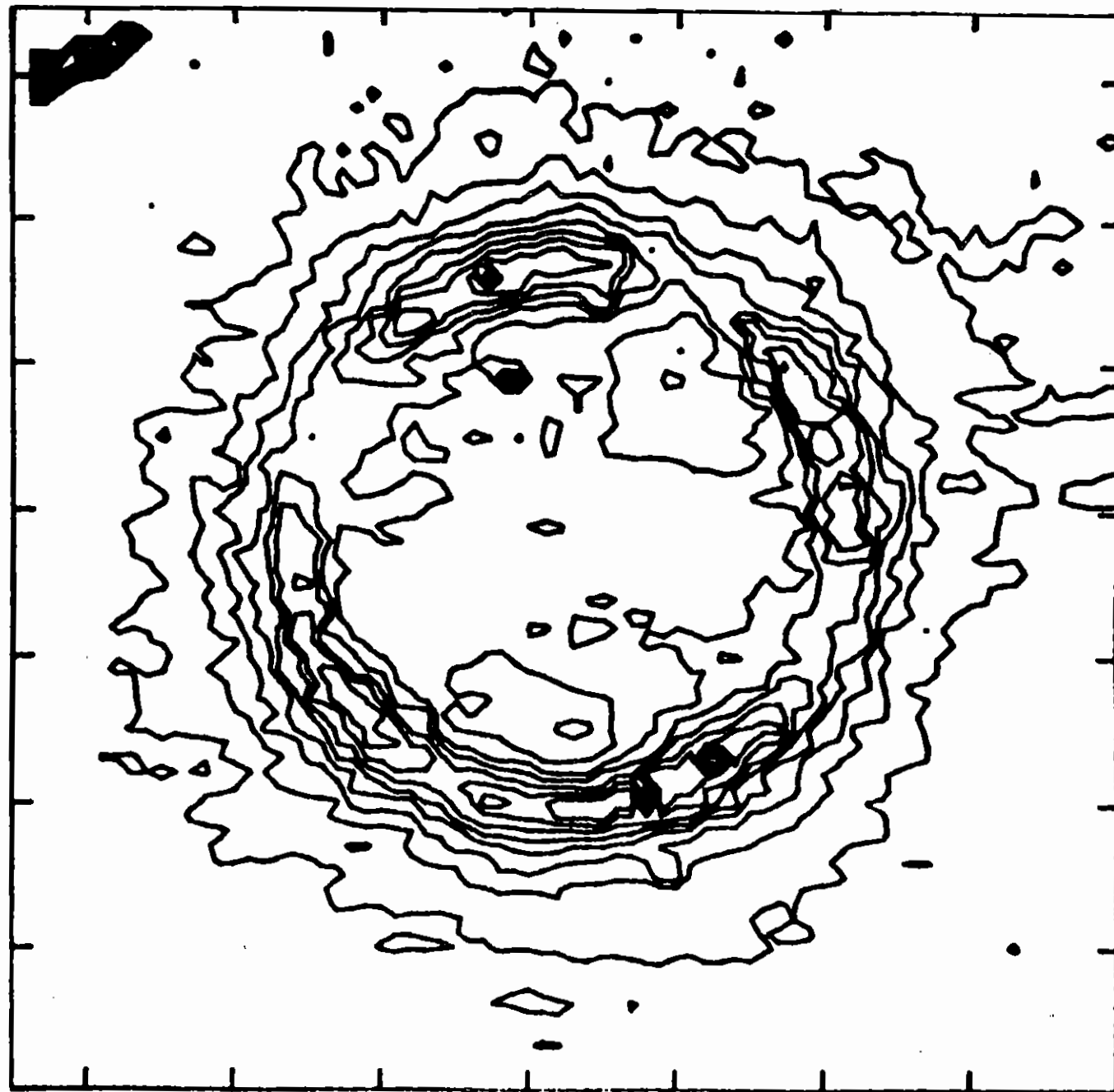


Fig. A2.3

Isodensity contour plot of X-ray image of solid glass target + $4\mu\text{m}$ of CH irradiated at $5 \times 10^{14} \text{ W cm}^{-2}$. The four peaks correspond to the four triple-beam overlap positions.

SHOT NUMBER = 9, ON 29/4/83

BEAM	IR ENERGY	GREEN ENERGY	EFFICIENCY

UE	91.59	41.44	0.56

UW	111.71	42.32	0.47

N	126.27	41.66	0.41

S	91.89	40.65	0.55

LE	112.00	40.12	0.44

LW	89.14	39.75	0.55

AVERAGE = 39.5% CONVERSION

TOTAL ENERGY = 245.93 JOULES, AVERAGE = 40.99 J/BEAM

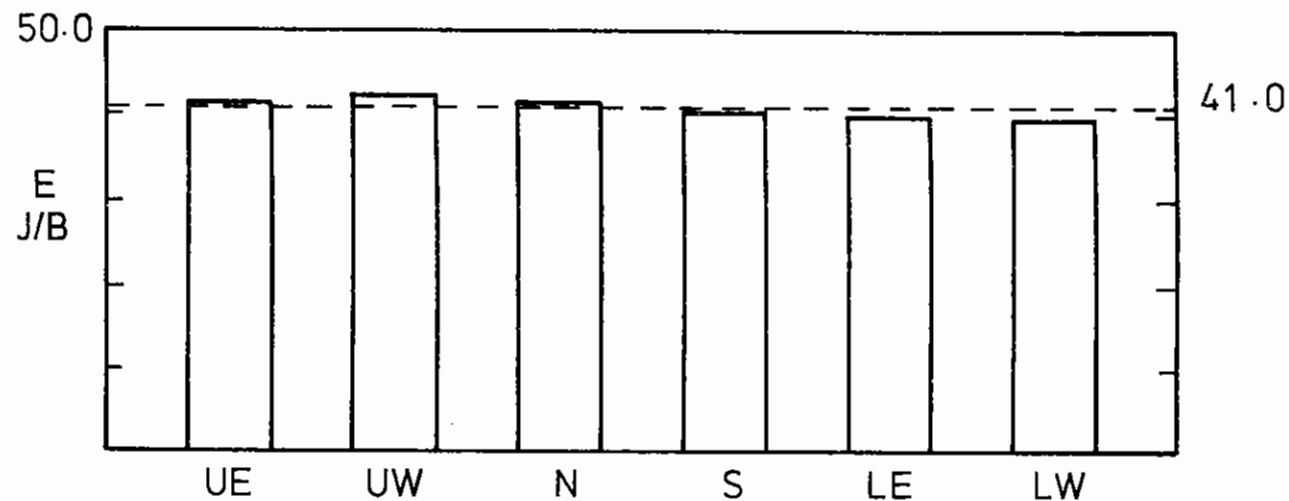


Fig. A2.4 Beam balance histogram associated with Fig. A2.3.

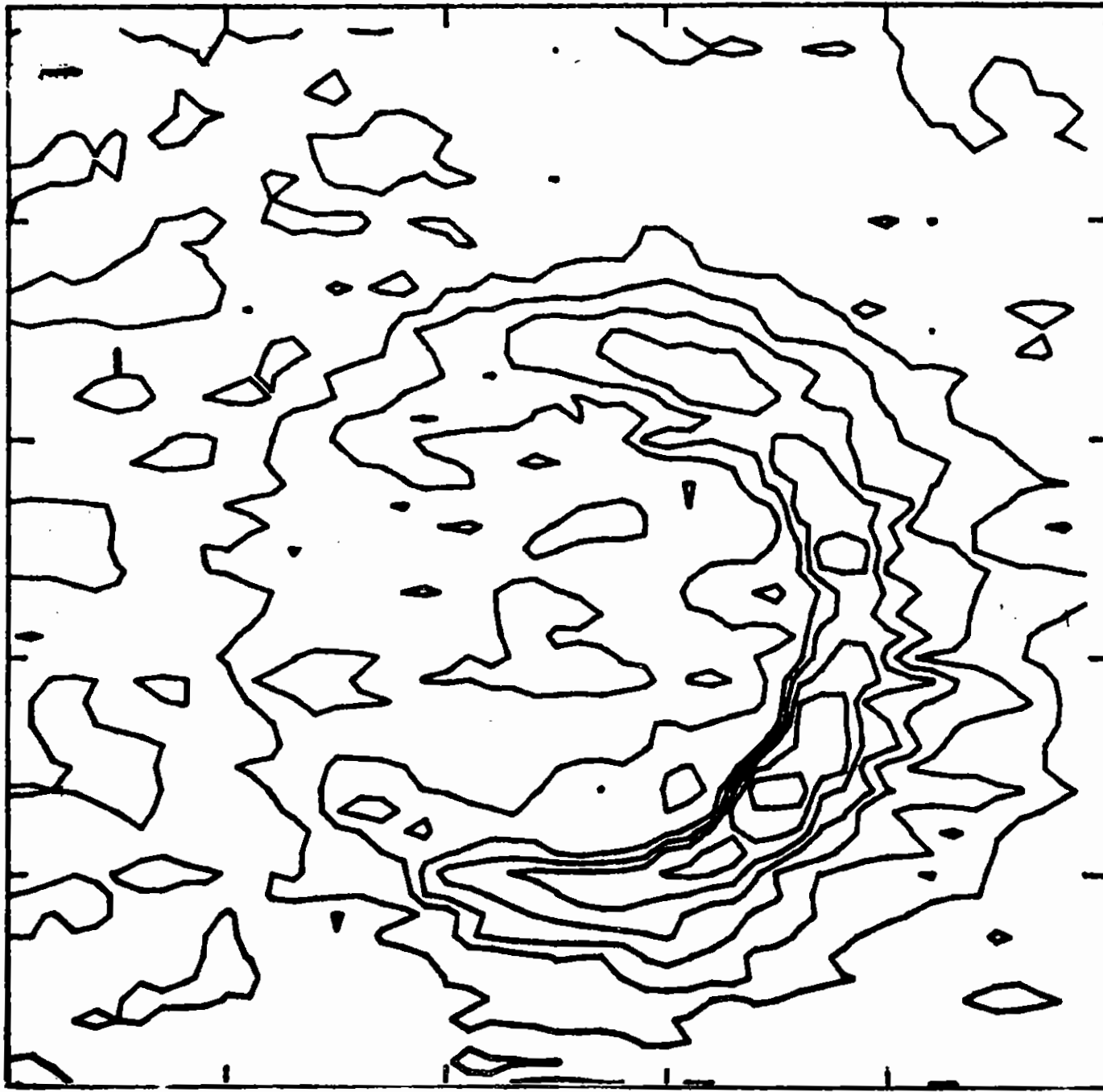


Fig. A2.5 Isodensity contour plot of target irradiated with poorly balanced beams.

SHOT NUMBER = 11, ON 29/4/83

BEAM	IR ENERGY	GREEN ENERGY	EFFICIENCY

UE	68.53	31.90	0.57

UW	84.13	28.08	0.41

N	95.56	35.62	0.46

S	69.22	29.43	0.52

LE	85.37	29.18	0.42

LW	66.85	28.40	0.52

AVERAGE CONVERSION = 38.6%

TOTAL ENERGY = 182.60 JOULES, AVERAGE = 30.43 J/BEAM

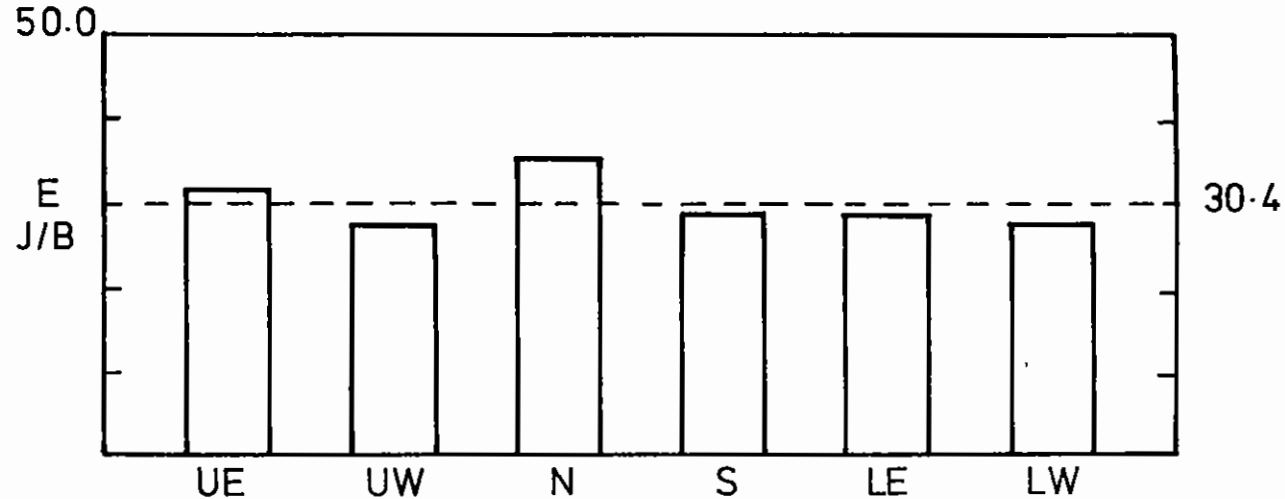


Fig. A2.6 Beam balance histogram associated with Fig. A2.5.

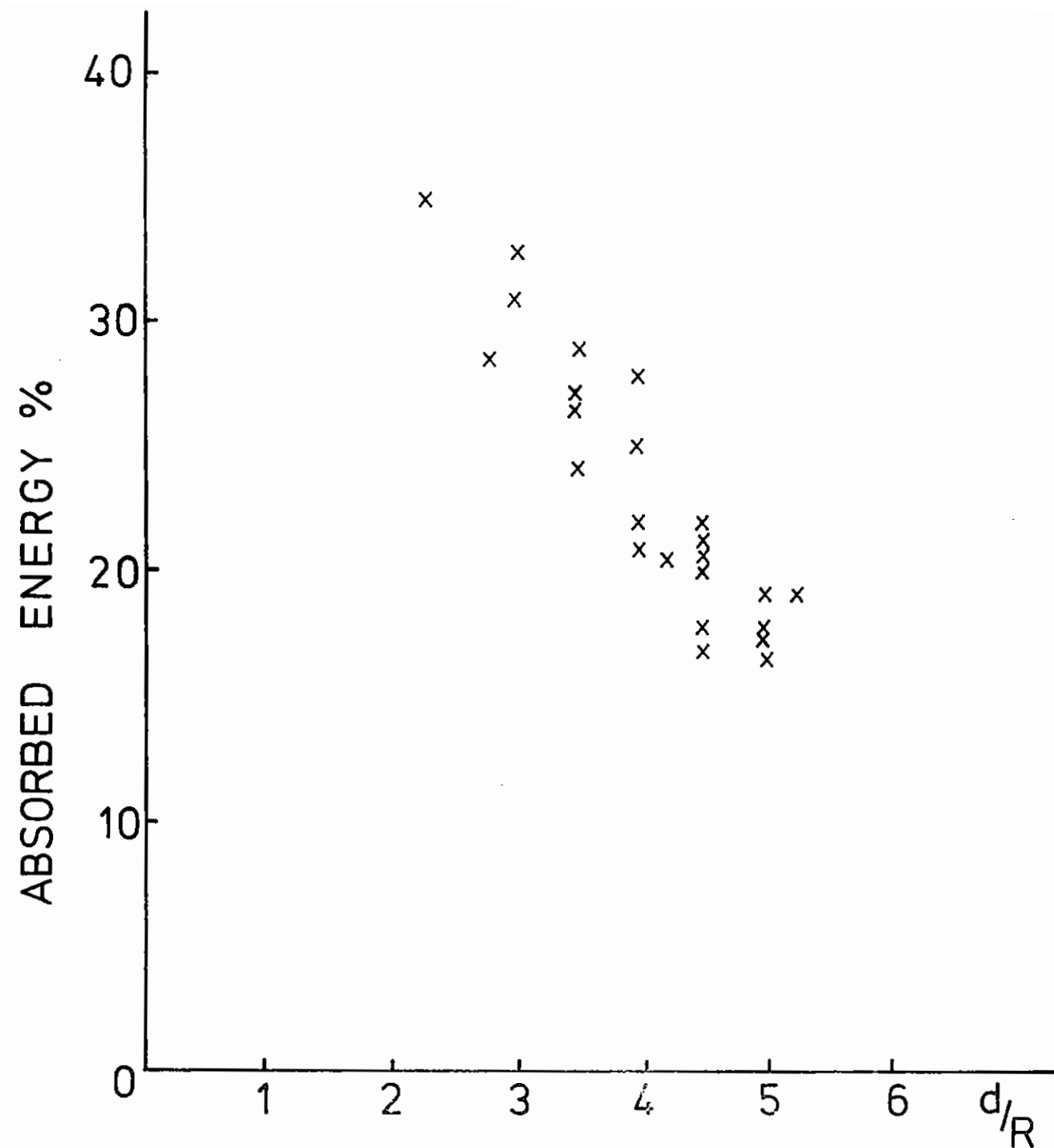


Fig. A2.7 Absorbed energy as a function of defocus parameter d/R .

probably due to refraction effects and the fact that the corona is larger than the initial target surface; even at a defocus of 5 the rays from the edge of the lens pass within one initial target radius of the target surface.

In conclusion we have observed a higher mass ablation rate at the triple beam overlaps, in qualitative agreement with a $\cos \theta$ or $\cos^2 \theta$ absorption dependence. However, such effects were often masked by poor beam balance which must be improved in further uniformity experiments. The absorbed energy has been measured as a function of defocus and shows that the calculated improvement in uniformity at defoci greater than obscuration (A2.7) does not produce as sharp a fall off in fractional absorbed energy as purely geometric considerations would suggest.

A2.2.3(b) Harmonic imaging studies of illumination uniformity for 0.53 μm ablative compressions.

V Aboites, S Sim (Essex), D Bassett and I Ross (RAL)

(i) Introduction

The production of energy through the implosion of D-T microballons by using laser beams requires a high degree of uniformity of the energy deposition in the surface of the target in order to achieve spherical convergence (A2.10).

In this experiment a study of the uniformity of illumination of imploded microballons was done when they were irradiated by the six beams of the RAL Vulcan laser under different focusing conditions.

Several linear and non-linear processes in the interaction of intense electromagnetic radiation with plasmas provide us with information about very specific regions of a plasma.

At the critical density region (n_{cr}), where $\omega_p = \omega_0$, (ω_0 refers to the frequency of the incident light on the plasma) resonance absorption can give

rise to electron plasma waves of frequency ω_0 which may couple either to electromagnetic waves of frequency ω_0 or to another electron plasma wave of frequency ω_0 to generate $2\omega_0$ radiation (A2.11 and A2.12).

Alternatively, generation of $2\omega_0$ also may be due to the Parametric Decay Instability (PDI) where near the critical density n_{cr} , an electromagnetic wave of frequency ω_0 decays into an ion acoustic wave and an electron plasma wave of frequency near ω_0 which by further recombination with electron plasma waves ω_0 or electromagnetic waves ω_0 may generate $2\omega_0$ radiation (A2.13).

Therefore, the critical density region of a plasma n_{cr} , can be identified with the region of generation of $2\omega_0$ radiation. The appearance of the $2\omega_0$ emission would be expected to be specular.

Similarly, generation of odd-integer half harmonic emission is due to processes occurring at or near the quarter critical ($n_{cr}/4$) region of the plasma, where $\omega_p = \omega_0/2$.

Two of these processes are the Two Plasmon Decay Instability (TPD) or Stimulated Raman Scattering (SRS), where the former corresponds to the decay of an incident electromagnetic wave into two electron plasma waves of frequency $\sim \omega_0/2$, and the latter corresponds to the decay of an incident electromagnetic wave into a scattered electromagnetic wave and an electron plasma wave both of frequency $\omega_0/2$ (A2.14 and A2.15)

In particular, electromagnetic waves of frequency $3\omega_0/2$ can be generated by the coupling of an electromagnetic wave of frequency ω_0 and an electron plasma wave produced either by TPD or SRS. The nature of the $3\omega_0/2$ light would be expected to be diffuse.

For this experiment, in order to study the smoothness of the critical (n_{cr}) and quarter critical ($n_{cr}/4$) surfaces of the plasma produced by different irradiation conditions, observations were made of the $2\omega_0$ and $3\omega_0/2$ radiation emitted by the imploded microballons.

(ii) Experiment

Glass microballons of diameter $85\mu\text{m}$ to $180\mu\text{m}$ coated with $4\mu\text{m}$ of plastic were irradiated at intensities up to 10^{15} Wcm^{-2} by using six green ($0.5\mu\text{m}$) laser beams. The pulse length was 1ns and the beams were focussed using $f/1$ lenses. The energy beam balance was $\sim 10\%$.

The collection of $2\omega_0$ and $3\omega_0/2$ emission was done by using an aluminium coated spherical mirror of 150mm focal length. An iris was used to stop the aperture of the mirror down to $f/10$ giving an experimental final resolution of approximately $5\mu\text{m}$ limited by spherical aberrations.

The $2\omega_0$ and $3\omega_0/2$ light was diverted by mirrors (HR @ $0.23\mu\text{m}$ and HR @ $0.35\mu\text{m}$ respectively) and imaged onto two open-shutter cameras.

An interference filter centered at 263nm plus one OX7 filter and two UG5 filters giving an attenuation of $\sim 10^{-15}$ for $0.5\mu\text{m}$ were placed in front of the $2\omega_0$ camera, and an interference filter centered at 351nm plus one OX9 filter and one UG1 filter giving an attenuation of $\sim 10^{-10}$ for $0.5\mu\text{m}$ were placed in front of the $3\omega_0/2$ camera.

The alignment of the spherical mirror was done by using a He-Ne laser injected through the $3\omega_0/2$ and $2\omega_0$ turning mirrors (see Figure A2.8). Once the spherical mirror and the microballon were aligned with the He-Ne laser, photographs of the static image of the microballon illuminated by the He-Ne laser were taken with both cameras. This was done in order to superimpose on the same film the image of the original position of the microballon with the image obtained by the $2\omega_0$ and $3\omega_0/2$ emission.

The focussing of the laser beams on the microballon was characterized by the parameter D/R where R is the radius of the microballon and D is the distance from the target to the focal plane of the focussing lens. Figure A2.9.

The collection aluminium mirror was placed equidistant from three of the final laser focussing lenses i.e. at an angle of $\sim 54^\circ$ from the axis of these three laser beams.

For our experimental conditions our computer simulations show that light which is specularly reflected by an ideal spherical target surface would be observed as shown in Figure A2.10 where the three symmetric internal spots correspond to light specularly reflected from the three front laser beams and the three symmetric external spots correspond to light specularly reflected from the three back laser beams. The dotted line represents the critical surface of the plasma.

Therefore, irregularities on the critical (n_{cr}) surface are manifest by observation of $2\omega_0$ radiation coming from regions of the target not corresponding to these six ideal spots.

(iii) Results

Figure A2.11 shows the result obtained from a surface tight focus shot, the three spots corresponds to the light emitted inside the angle of our collection optics. Figures A2.12 and A2.13 for a $D/R = 3.5$ and $D/R = 4$ respectively, show evidence of an irregular critical and quarter critical surface due to non uniform illumination. Finally, Figure A2.14 shows that for $D/R = 5$ we obtain the closest resemblance to the calculated image by our computer simulations for the case of specular light from an ideal spherical target surface Figure A2.10, which indicate that for $D/R = 5$ we get the smoothest critical surface for our experimental conditions. It should be noted that the position of the spot near the stalk of the microballon is lower than expected, this may be due to the disturbance of the plasma symmetry of the microballon produced by the presence of the expanded plasma from the stalk. This is corroborated by the $3\omega_0/2$ emission picture which shows a pear shaped quarter critical surface.

Our results show, in agreement with our computer simulations, that for our experimental conditions the best degree of uniformity of illumination is obtained for a parameter $D/R = 5$, when account is taken of the stand-off of the critical surface from the solid and the need to avoid regimes in which the uniformity is particularly poor.

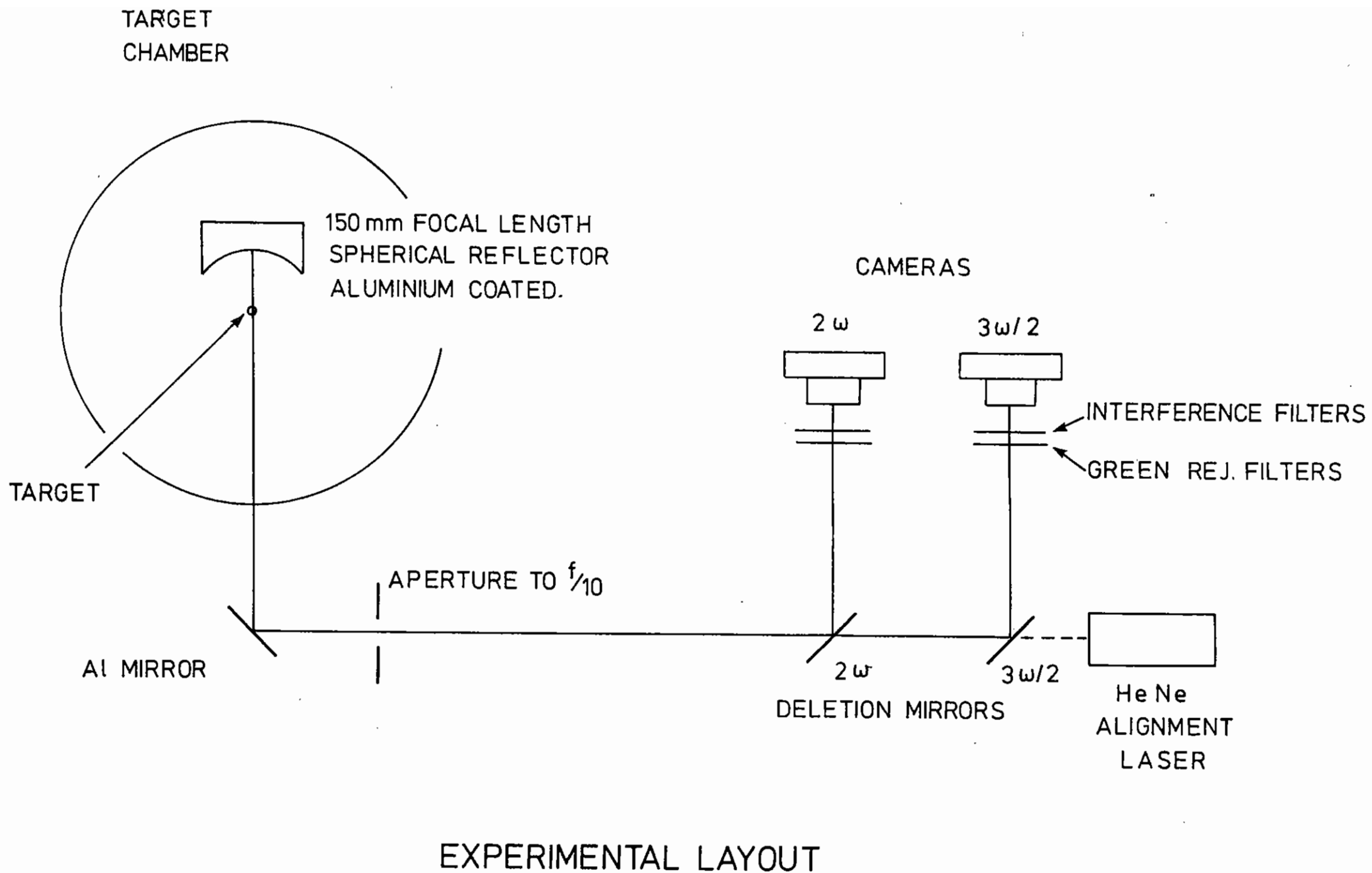
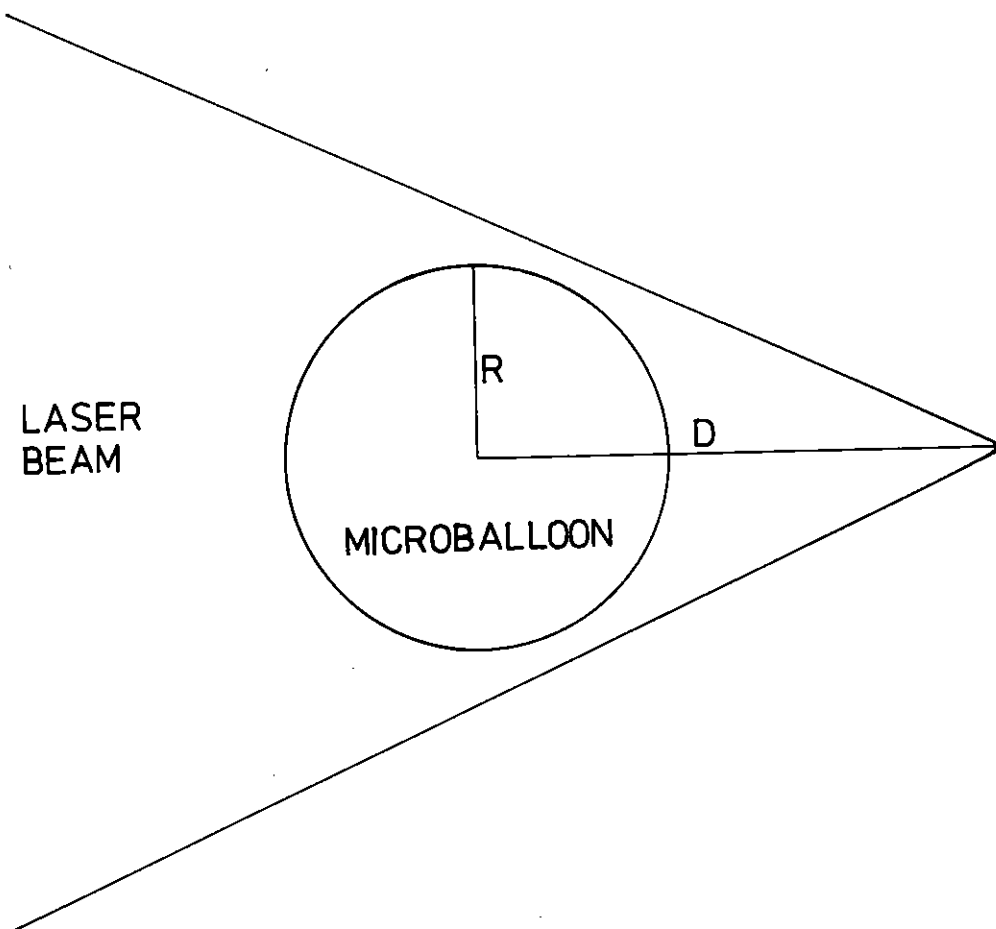


Fig. A2.8 Experimental arrangement for harmonic imaging studies at $3\omega_0/2$ and $2\omega_0$.



FOCUSING PARAMETER = D/R

Fig. A2.9 Schematic diagram of defocus conditions.

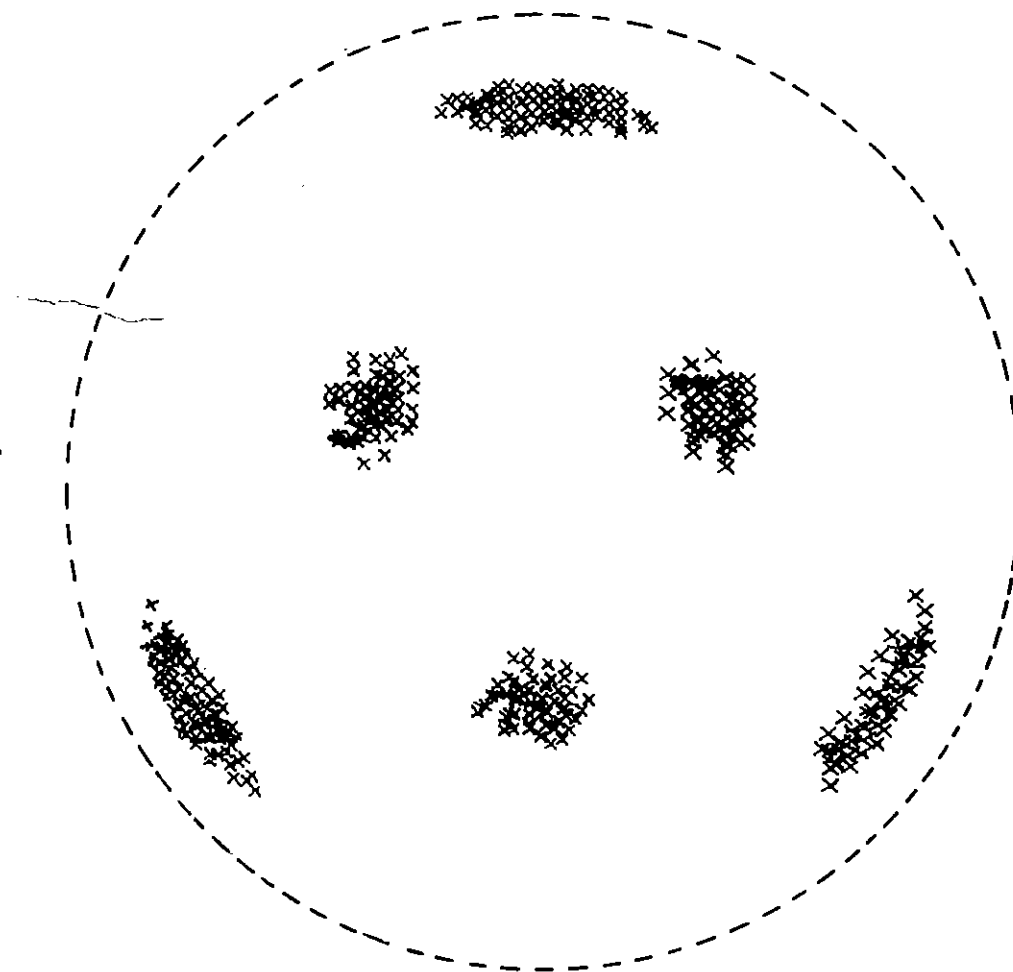


Fig. A2.10 Computer simulation of expected specular emission regions of a microballoon target.

SURFACE TIGHT FOCUS

2ω
EMISSION



$\frac{3\omega}{2}$
EMISSION



Fig. A2.11 Harmonic images for a tight surface focus.

$$\frac{D}{R} = 3.5$$

2ω
EMISSION



$\frac{3\omega}{2}$
EMISSION



Fig. A2.12 Harmonic images for a defocus of $D/R = 3.5$.

$$\frac{D}{R} = 4$$

2ω .
EMISSION



$\frac{3\omega}{2}$.
EMISSION

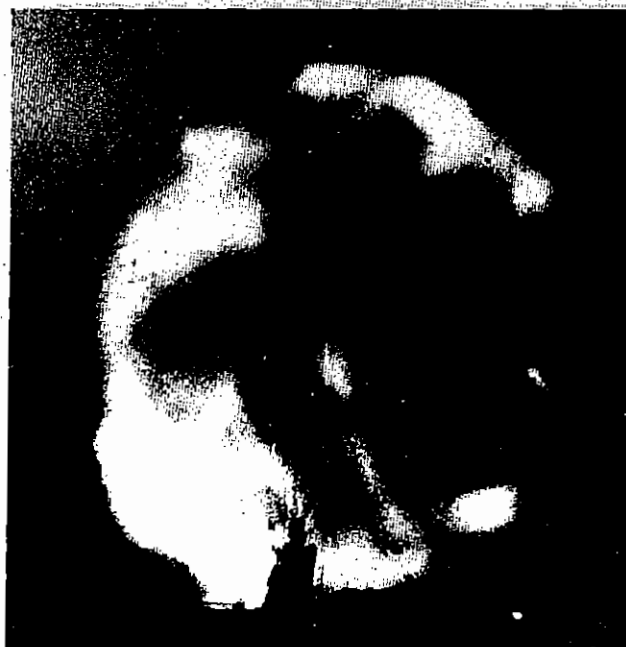


Fig. A2.13 Harmonic images for a defocus of $D/R = 4$.

$$\frac{D}{R} = 5$$

2ω .
EMISSION



$\frac{3\omega}{2}$.
EMISSION

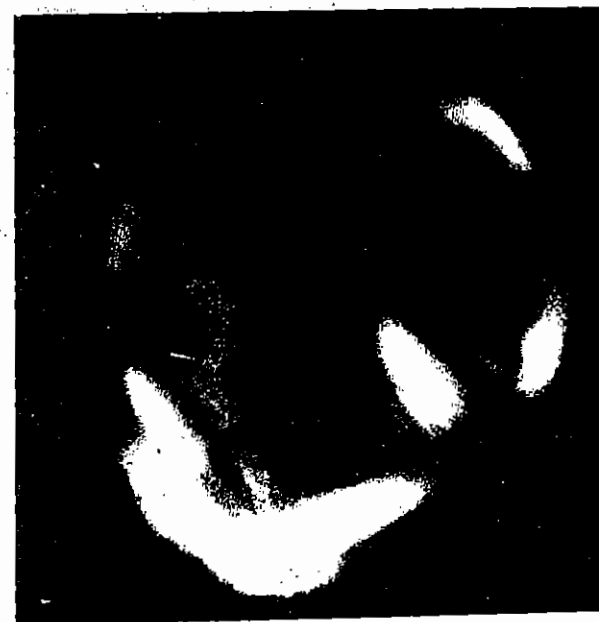


Fig. A2.14 Harmonic images for a defocus of $D/R = 5$.

In order to ensure that the uniformity or non-uniformity of the emission observed was effectively due to the D/R parameter and not to lack of uniformity in the Laser beam profiles, pictures were taken displaying the beam profiles for all six beams for every shot. Figure A2.15 shows the $2\omega_0$ and $3\omega_0/2$ emission together with two typical samples of the beam profiles for the same shot.

It was found that non-uniform laser beam profiles were obtained whenever the laser is fired at a repetition rate faster than approximately 45min.

It can be seen in Figure A2.15(a) that for a qualitatively uniform laser beam profile we get a very non uniform critical (n_{cr}) surface which is due to the focussing condition $D/R = 2.25$ used in this shot.

For a focussing condition of $D/R = 5$, which is the optimum condition for uniform illumination, we can see in Figure A2.15(b) the lack of smoothness of the critical surface caused by the very poor uniformity of the laser beam profiles. The reason of this poor laser beam profile is that this shot was fired just 20 minutes after the previous one.

However, Figure A2.15(c) show that for the same focussing condition $D/R = 5$ and uniform laser beam profiles obtained because the laser was fired 78 minutes after the previous shot, we obtain a pattern indicating a smooth critical surface.

(iv) Conclusions

From experiments of imploded microballons with six laser beams it was found that the optimal condition for uniformity of illumination is at the the focussing condition $D/R = 5$. This value is in agreement with computer simulations which predict optimal illumination for $D/R > 4.5$.

The importance of uniform laser beam profiles has also been established.

Future experiments on the new twelve beam laser are planned.

A2.15

- $2\omega_0$ - - $3\omega_0/2$ - - BEAM PROFILE -

a)

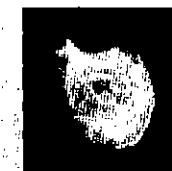
$$\frac{d}{R} = 2.25$$

SHOT 07/220483
TARGET DIAM 266 μ m
COATED WITH 200nm Au
ENERGY ON TARGET 211J.

- $2\omega_0$ -

- $3\omega_0/2$ -

- BEAM PROFILE -



b)

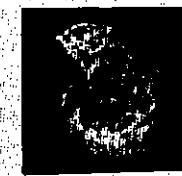
$$\frac{d}{R} = 5$$

SHOT 14/290483
TARGET DIAM 149 μ m
COATED WITH 4.6 μ m PLASTIC
ENERGY ON TARGET 199J

- $2\omega_0$ -

- $3\omega_0/2$ -

- BEAM PROFILE -



c)

$$\frac{d}{R} = 5$$

SHOT 10/270483
TARGET DIAM 137 μ m
COATED WITH 4 μ m PLASTIC
ENERGY ON TARGET 160J

- $2\omega_0$ -

- $3\omega_0/2$ -

- BEAM PROFILE -



Fig. A2.15

harmonic images showing the effect of beam profile on the observed uniformity.

A2.2.4 X-Radiographic Studies of Ablatively Imploded Plastic Shells

C L S Lewis, S Saadat, M J Lamb, P McCavana, J McGlinchey (QUB), M H Key,
W T Toner and R W Eason (RAL)

A2.2.4(a) Introduction

The ability to generate high density, hot plasma (10's of g/cm^{-3} and 100's of eV) by means of an ablatively imploded target depends largely on the degree of volumetric compression which can be achieved in practice and thus the degree of spherical symmetry and stability which can be maintained during the target convergence. There are several processes involved in the laser-target interaction and the hydrodynamic response of the target which will influence the overall performance of the target and the quality of the implosion attained.

Potential problems originate with the deposition of laser energy outside the critical surface and its subsequent transport into the ablation surface. With a finite number of beams incident on the target the uniformity of the absorbed energy density in the critical region depends on several factors. These include the number of beams, the energy balance between beams, the small scale intensity fluctuations within a beam and the overlapping of beams in the vicinity of the target. The extent of overlapping is related to the focussing lens effective f/number, the beam profiles, the size of the target and its position relative to the lens focus. In general spatial variations in absorbed energy density are to be expected with a wide range of spatial frequencies. If these non-uniformities are mapped onto the ablation surface with significant amplitude variation then a non-spherical implosion will be driven. The exact nature of the mapping, i.e. the 2-D non-linear heat flow between the critical and ablation surfaces is not fully understood but some thermal smoothing is expected which will effectively smooth the higher spatial frequencies. The hydrodynamic response of the target will aid the smoothing process with an overall effect sensitive to the average irradiance on target and the laser wavelength. A counteracting effect resulting from the non-uniformities may be the generation of magnetic fields in the overdense plasma leading to inhibition of the thermal

smoothing process and possibly the generation of filamentary structure in the overdense plasma.

The integrity of the target is subject to the Rayleigh-Taylor instability during the initial acceleration phase (and later during the deceleration phase) but it is not clear how the growth of this hydrodynamic instability may be damped by the above thermal and hydrodynamic smoothing processes.

Decompression of the target shell during the implosion will limit the volumetric compression and peak density attained. Decompression, as opposed to shell distortion through non-uniform drive pressure or Rayleigh-Taylor break-up, may result from preheating through shocks, soft X-ray radiation from the surface plasma or hot electrons. Many of the processes or limited combinations of them can be studied in isolation per se and such experiments are reported elsewhere in this report. It was the objective of this experiment to identify on a reasoned but somewhat empirical basis a target design and a laser illumination prescription which minimised the adverse effects and produced an implosion which could be adequately described by the best modelling available. Having obtained a numerical model 'fix', future experiments could be designed to investigate the effects of the above processes by stretching experimental parameters to the point that present modelling breaks down.

A2.2.4(b) Experimental arrangement

The six beams of the VULCAN laser were used in second harmonic to deliver $\lesssim 150$ J of incident energy in 0.9 nsec to low aspect ratio, unfilled polystyrene shell targets. The principal diagnostic was pulsed X-ray radiography using pinhole imaging and an active phosphor/image intensifier detection system. The experimental set-up is illustrated in Figure A2.16 and the pinhole camera diagnostic axis relative to single, double and triple beam overlap directions is illustrated in Figure A2.17. The camera axis is along a triple beam overlap direction.

Targets with a typical wall thickness of $\Delta R \approx 10 \mu\text{m}$ were chosen to minimise decompression effects through radiative preheat (A2.3, A2.16) (See also

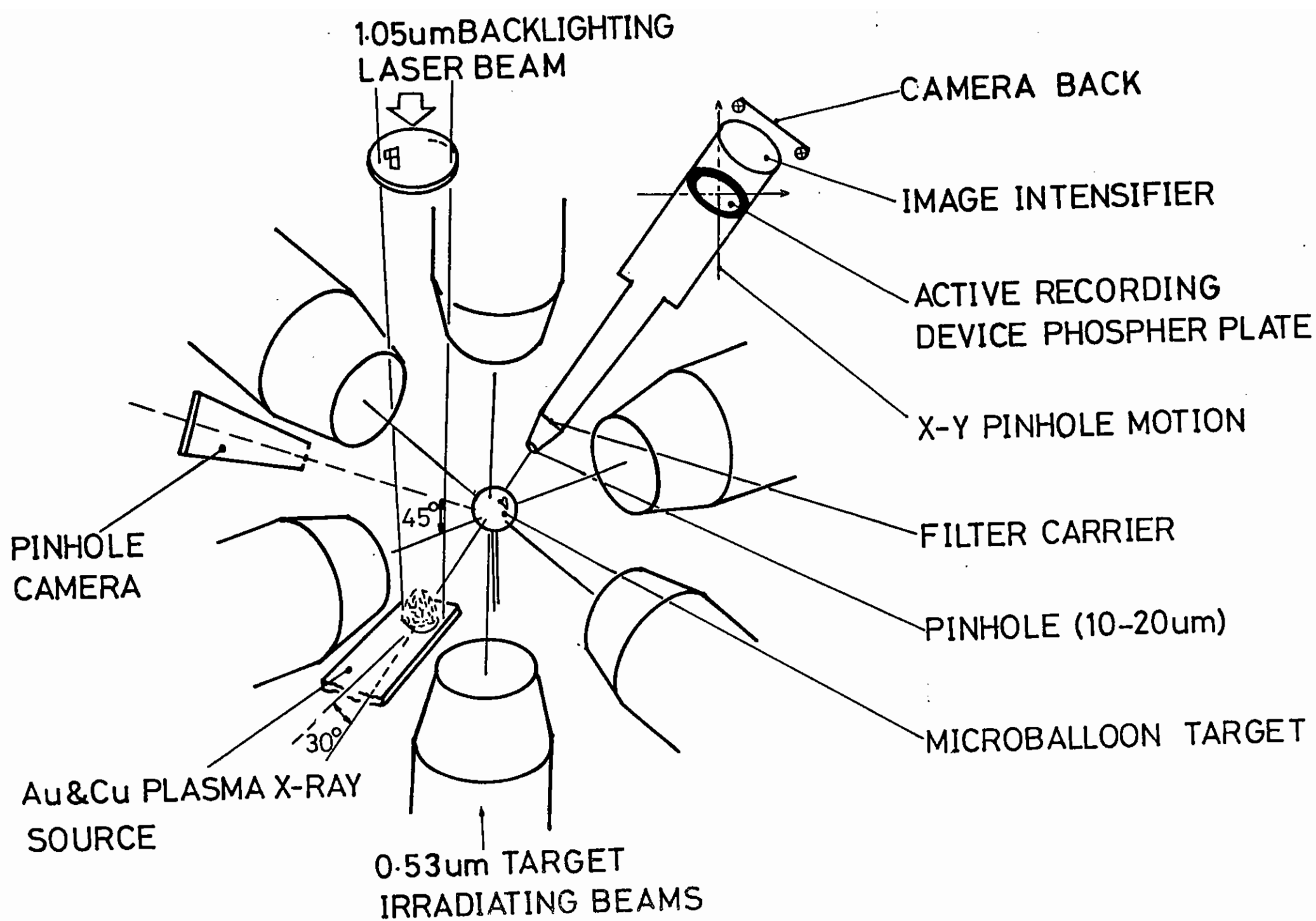


Fig. A2.16 Experimental arrangement.

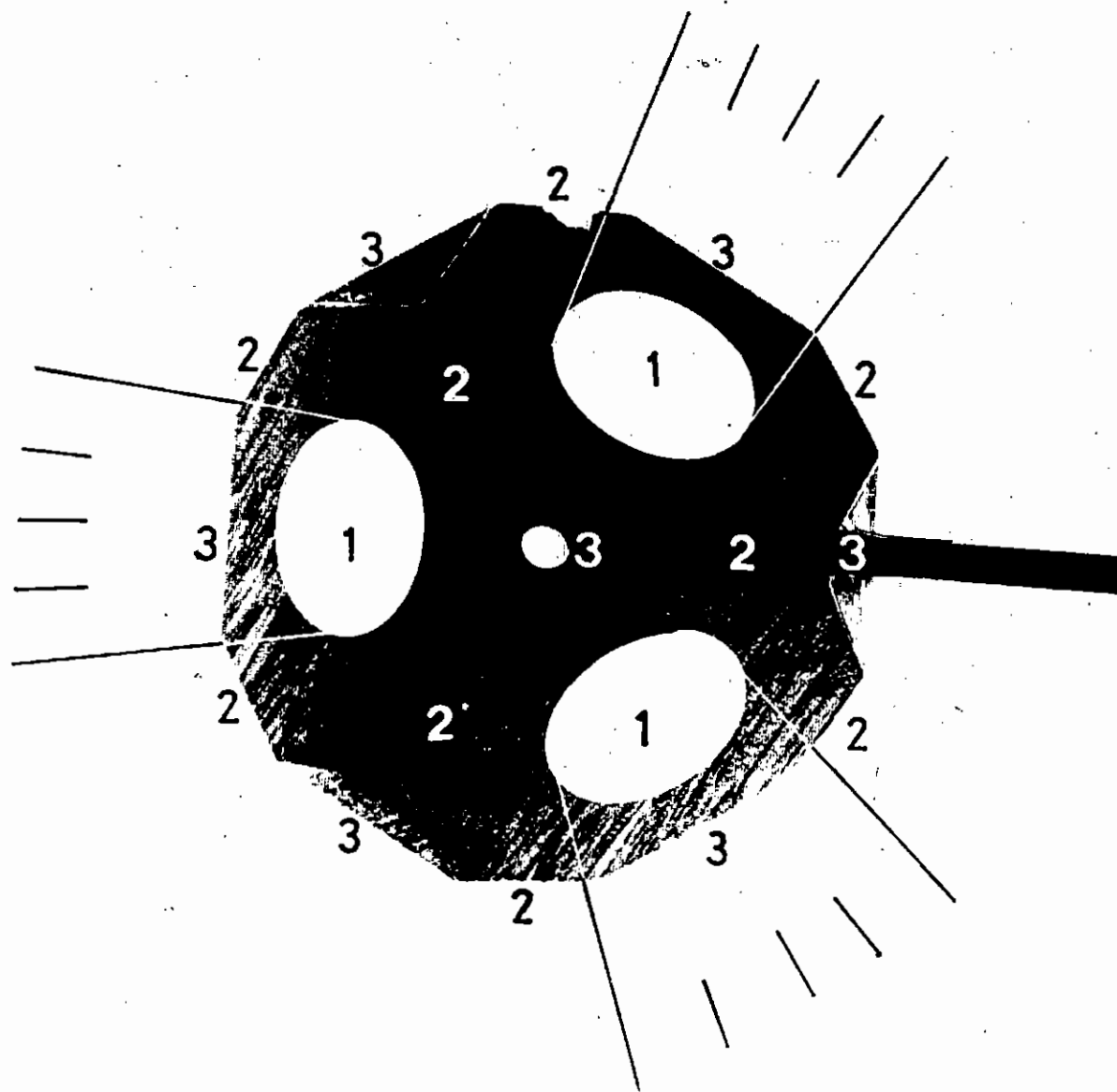


Fig. A2.17

Beam geometry with respect to viewing axis. Symbols 1, 2 and 3 refer to centre of beams, two-beam overlaps and 3-beam overlaps respectively.

section A5.4.6 *ibid*). Shell diameters were typically $2R \sim 160 \mu\text{m}$ to achieve ≤ 5 MBar drive pressure with the final irradiance conditions and to simultaneously give a relatively low aspect ratio $R/\Delta R \sim 8$ to minimise any Rayleigh Taylor instability and allow for a moderately high final core density. The implosion time scale of 1-2 nsec was sufficiently long to allow smoothing processes to approach steady state (A2.5).

Uniformity of target illumination was approached empirically (as described in detail in Section A2.2.3 *ibid*) by observing the uniformity of X-ray, second harmonic and three halves harmonic emissions from the target as a function of the defocussing of the main lenses. Uniformity increased with D/R (where D is the distance of the focal plane beyond the target centre and R is the target radius) but at the expense of energy absorbed into the target which was measured with plasma calorimeters. A compromise was adopted with D/R = 4.5 and a fractional laser energy absorption of 0.2. Typical absorbed irradiance levels were then $1-5 \times 10^{13} \text{ W cm}^{-2}$. Beam energy calorimetry showed that the six beams had $\sim 3\%$ rms energy deviations from the mean energy per beam and target equivalent plane monitors showed each beam had typically $\pm 20\%$ variations in intensity about the beam average on a spatial scale of $\sim 30 \mu\text{m}$ on the target surface.

A2.2.4(c) Diagnostics

(i) X-ray backlighting

$\sim 30 \text{ J}$ of $1 \mu\text{m}$ radiation in a 100 ps pulse was focussed to $\sim 300 \mu\text{m}$ spot size on a backlighting target (Cu or Au foils angled at $\sim 30^\circ$ to the pinhole camera viewing direction) and was driven by an oscillator synchronised to a long pulse oscillator driving the main target beams. The backlighting plasma was $\sim 1 \text{ mm}$ from the main target. The imaging pinhole was positioned $\sim 7 \text{ mm}$ from the main target and a P 11 phosphor coated fibre-optic plate detector doubled as a chamber window $\sim 50 \text{ cm}$ away. A fibre-optically coupled Mullard 50/40 image intensifier outside the chamber recorded optical images on HP5 film with a demagnification of ~ 0.7 . Pinholes were fabricated with an off-line Q-switched YAG laser in 15μ thick Pt foil and referenced spatially in a mounting stud so that pinholes could be replaced

from shot-to-shot with a $10 \mu\text{m}$ accuracy. $12 \mu\text{m}$ diameter pinholes were used with the copper backlighting targets and a permanent $15 \mu\text{m}$ Be filter for early time soft X-ray radiography. $20 \mu\text{m}$ diameter pinholes with an additional $30 \mu\text{m}$ saran ($\text{C}_2\text{H}_2\text{Cl}_2$) filtering were used with gold backlighting targets for later-time harder X-ray radiography.

(ii) X-ray streak camera

This was used to streak the self emissions of a typical shell target and a typical backlighting target on the same shot. X-rays projected through a 1 mm wide slit covered with a filter pack providing a range of cut-off energy channels from 1.5-4.0 KeV were streaked with a time resolution of $\sim 30 \text{ psec}$ to measure the relative timings of the target and backlighting beams and also to give a measure of the frame time of a radiograph. This was typically $\sim 200 \text{ psec}$ for a copper target as illustrated in Figure A2.18 and $\sim 100 \text{ psec}$ for a gold target.

(iii) Self-emission imaging

A $\times 10.7$ magnification pinhole camera with Kodak No-Screen film detection was used to image both shell and backlighting self-emission to keep a record of the shell illumination uniformity and the relative self-emission brightnesses. Two channels were used employing $5 \mu\text{m}$ and $25 \mu\text{m}$ Be. In all data shots the CH shell self-emission was negligibly small.

(iv) Backlighting spectral content

Miniature crystal spectrometers with Tl AP and PET crystals and SB5 X-ray film were used to monitor the Cu and Au backlighting spectra respectively. The relative energy content of the backlighting sources when filtered for the appropriate pinhole imaging channel was deduced and is shown in Figures A2.19 and A2.20 for Cu and Au respectively. An estimate of the absolute brightness was possible in the latter case but not needed for the data reduction later.

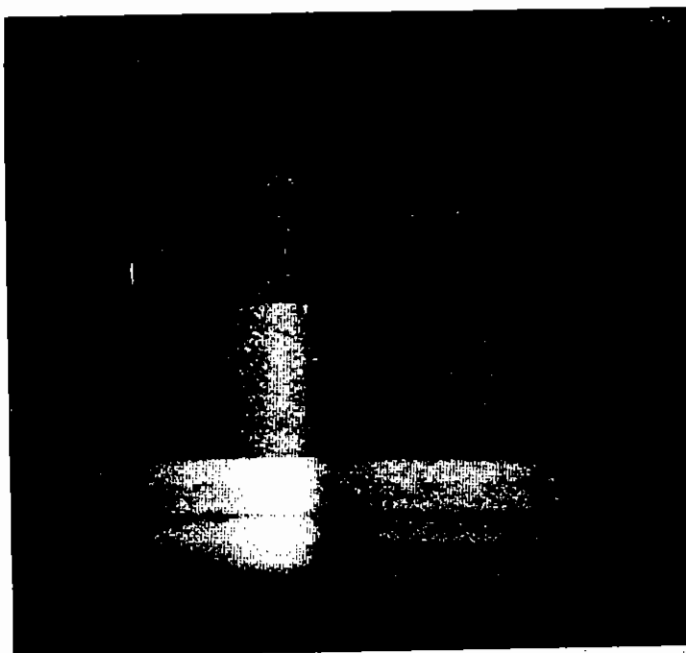
X-RAY STREAK OF THE TWO LASER PULSES

4.15 KEV

3.8 KEV

3.3 KEV

1.65 KEV



INTENSITY

10

1

3.3 KEV

T

1 ns

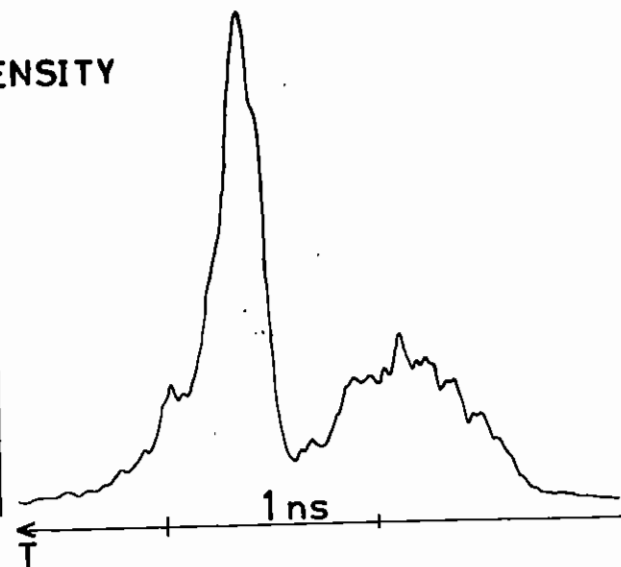
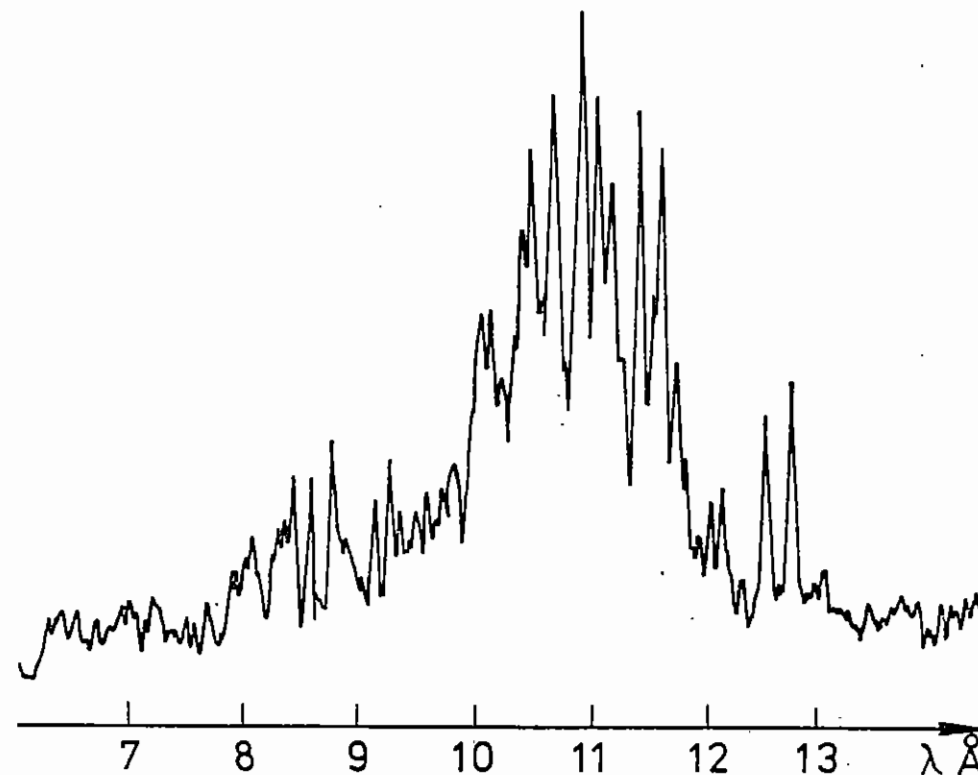


Fig. A2.18

Self-emission streak of CH shell and Cu backlighter in energy channels 1.6 - 4.2 keV.

Cu SPECTRUM



Rel. Int.

5

0

1.0

1.5

2.0

E keV

15 μm Be

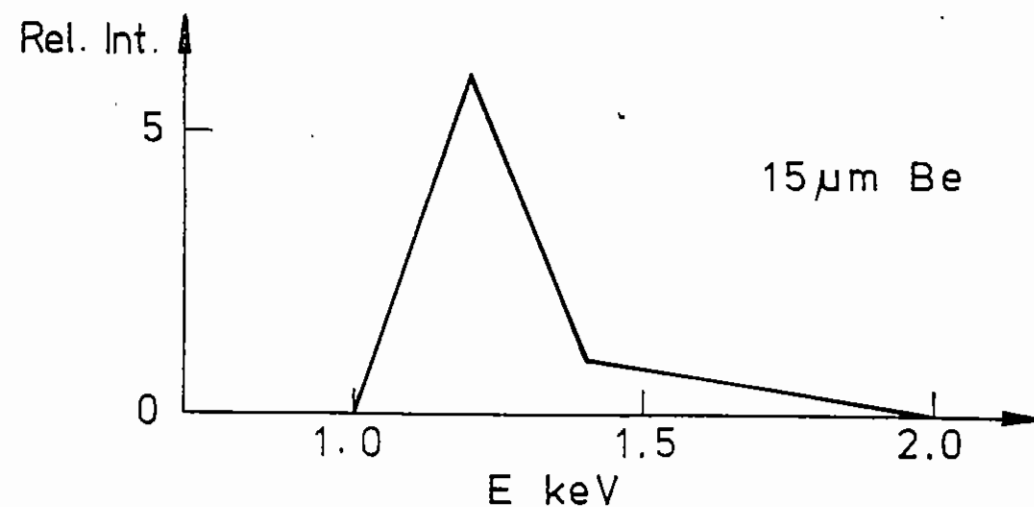


Fig. A2.19

Relative energy content of copper backlighting spectrum.

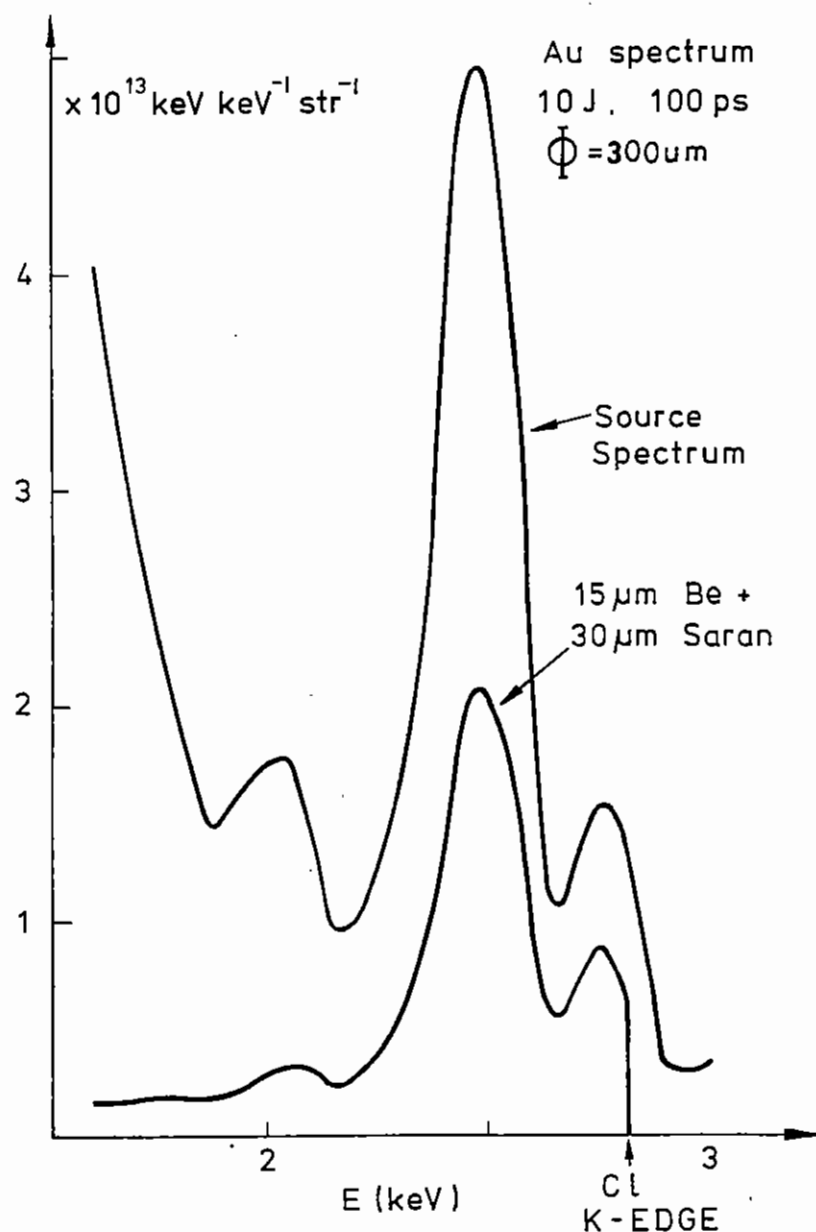


Fig. A2.20 Absolute energy content of gold backlighting spectrum.

A2.2.4(d) Radiography Data

Raw radiographs are shown in Figures A2.21 to A2.30 and experimental parameters listed in Table A2.2. Figure A2.21 is a system calibration shot where an unirradiated CH shell has been radiographed with a Cu source ($h \sim 1.2 \text{ keV}$) to provide the overall magnification ($\times 50$) and the overall spatial resolution in the target plane ($\sim 15 \mu\text{m}$). The fading of backlighting image intensity in the upwards direction within the shell corresponds to a softer X-ray spectral content in the backlighting plasma plume.

Figures A2.22 to A2.26 show Cu radiographs taken at early stages of convergence while Figure A2.27 is close to peak implosion but shows no measurable transmitted light through the core feature. Figures A2.28 and A2.29 show Au radiographs ($h \sim 2.5 \text{ keV}$) near peak compression with sufficient transmitted intensity to estimate the core density. Figures A2.22 to A2.30 show the initial shell dimensions to scale but centred with respect to the radiograph by eye as no absolute reference was available. Figure A2.30 shows a misaligned radiograph but illustrates the geometric centring of the core with respect to the shell self emission which is also seen to be negligibly weak compared to the backlight.

Several qualitative features are clear from the raw data. Early radiographs, delineating the shell outer boundary, show hexagonal structure reflecting the six beam geometry used and the associated spatial variations in absorbed energy density. This structure is to be expected as illustrated in Figure A2.31 where a simple model prediction of the target shape is compared to a data shot. The model assumes an absorption coefficient scaling as $\cos^2\theta$ (θ the angle of incidence) and a fourth power supergaussian beam profile. The shape, as seen by the pinhole camera, is calculated when the initial radius r_0 is reduced to $r_0/3$ for a lens defocus of $D/R = 4.5$. The distortion from spherical symmetry is due to $\sim 15\%$ higher drive pressure along a triple beam overlap direction compared to a single beam direction. At late times Figures A2.27 to A2.29 this structure seems to have largely smoothed out, to within the spatial resolution of the data.

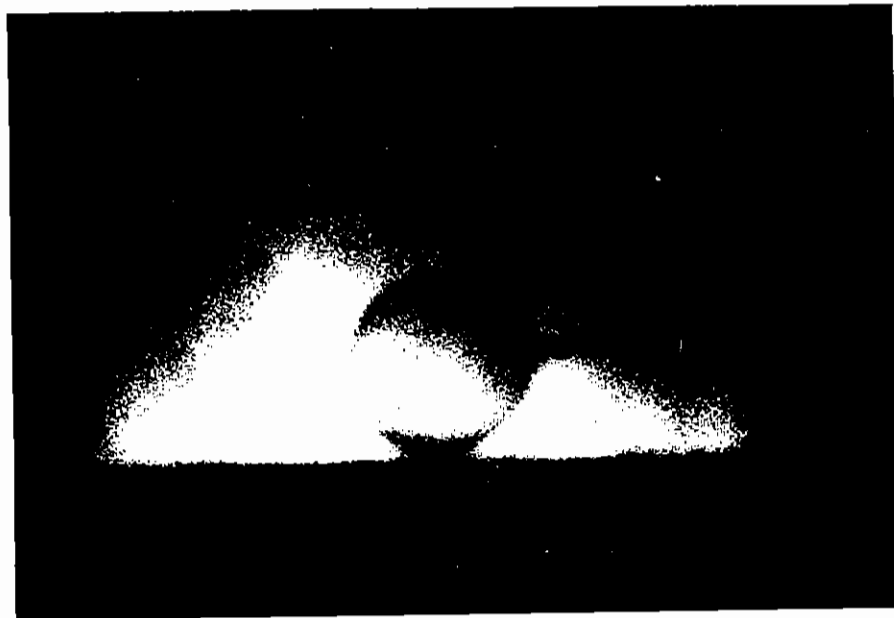
Date	Shot No	Target	ϕ μm	ΔR μm	$E_{2\omega\text{in}}$ J	$E_{\text{BIL inc}}$ J	$I_{2\omega\text{inc}}$ 10^{14} W/cm^2	I_{abs} 10^{13} W/cm^2	τ probe* PS	D/R	r_1 exp μm	r_1 med μm	Peak density g/cm ³ exp. Able inversion	τ_{imp} ns	Med peak density g/cm ³	T probe ns	B/L Target	Figure A2.35 Label
19/5/83	19	3413	159.8	11	129.9	20.4	1.8	3.6	600	4.5	45-57.5	43.9		2.3	3.5	1.680	Cu	1
19/5/83	29	3409	162.03	11.25	90.63	22.6	1.22	1.83	700	5.5	46.5	54.4		3.0	2.56	1.780	Cu	2
20/5/83	01	3441	147.46	9.78	171.2	25.2	2.78	4.17	700	5.5	36.6-46.4	28.3		2.0	6.0	1.780	Cu	3
20/5/83	04	3437	189.04	13.5	132.055	19.3	1.31	2.62	600	4.5	62.5-64	65.6		3.2	2.5	1.680	Cu	4
20/5/83	09	3442	150.35	9.9	166.94	21.3	2.61	5.50	500	3.5	36	29.1		1.87	4.9	1.580	Cu	5
26/5/83	23	3418	149.62	7.9	134.4	24.1	2.21	4.25	700	4.5	30	19.7		1.9	58	1.780	Cu	6
27/5/83	6	3422	167.17	8.3	101.1	24.2	1.28	2.56	700	4.5	51.7-57.8	40.8		2.1	3	1.780	Cu	7
8/6/83	7	3451	150.38	8.32	124.4	27	1.94	4	700	4.5	34-35.8	23.9		1.95	10	1.780	Cu	8
8/6/83	8	3440	178	12.8	76.23	22.7	.85	1.7	800	4.5	64.5 $\pm$.5	64.3		3.58	2.0	1.880	Cu	9
9/6/83	1	3446	163.52	11.5	91.8	24	1.31	2.4	1100	4.5	35	34.8		2.76	3.2	2.180	Au	10
9/6/83	11	3453	156	9.2	135.3	27.4	1.96	4	800	4.5	26	23.2		2.1	14	1.780	Au	11
10/6/83	1	3455	149	7.2	103.6	33.8	1.65	3.3	650	4.5	18.6	19.0	10	1.9	8.0	1.730	Au	12
10/6/83	3	3457	144.5	7	115.2	33.1	1.95	3.9	550	4.5	20.4	20.4	9	1.8	8.2	1.630	Au	13

*Probe timed with respect to peak of drive pulse

TABLE A2.2

SHOT 29 ON 19/05/83

X-RAY SHADOW OF A COLD CH SHELL



x47
 $\Delta X = 10 \mu\text{m}$

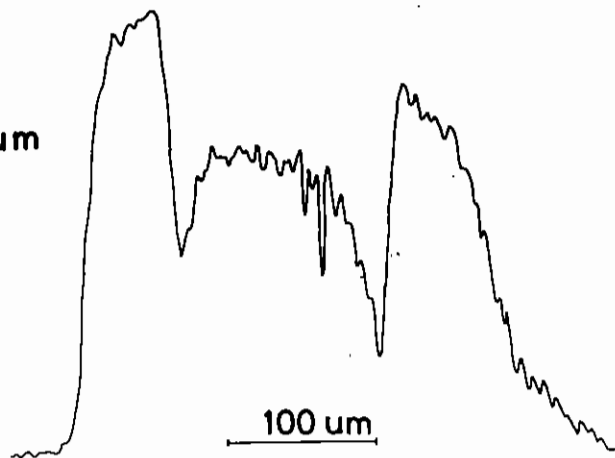


Fig. A2.21 Radiograph of unirradiated CH shell.

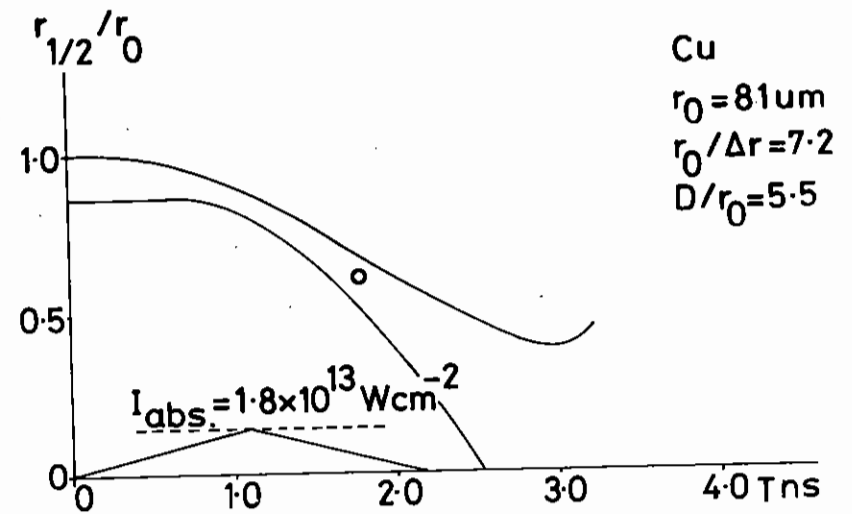
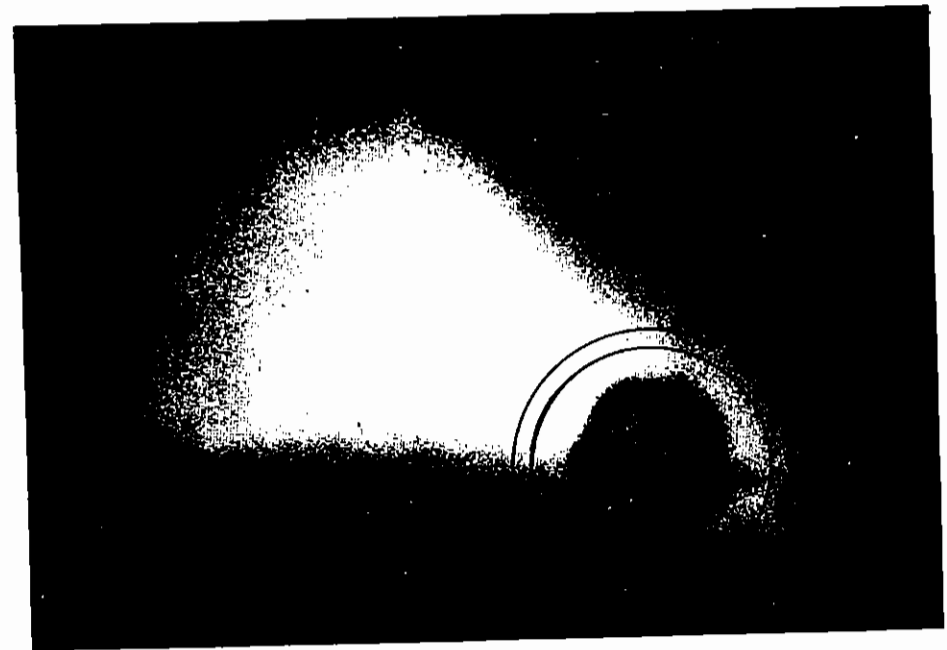


Fig. A2.22 Copper spectrum radiograph of irradiated CH shell.

SHOT 06 ON 27/05/83

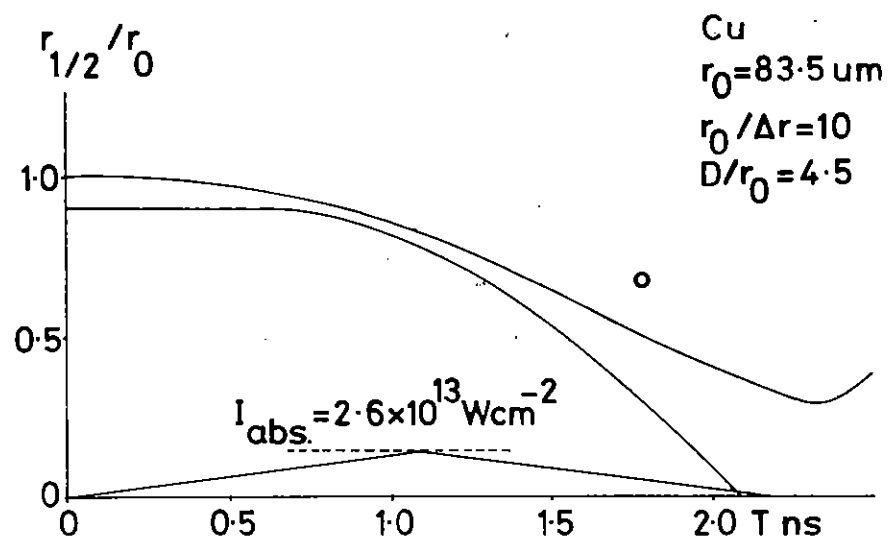
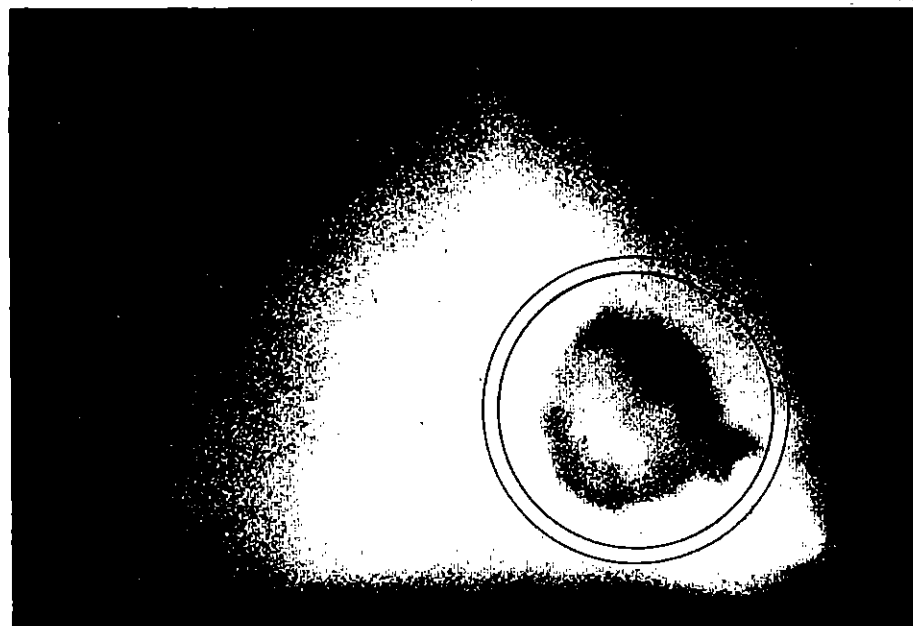


Fig. A2.23 Copper spectrum radiograph of irradiated CH shell.

SHOT 04 ON 20/05/83

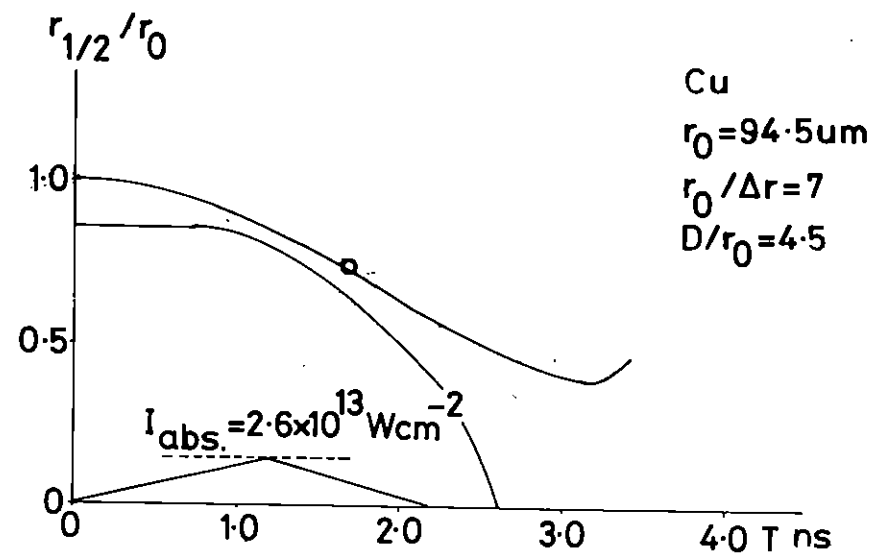
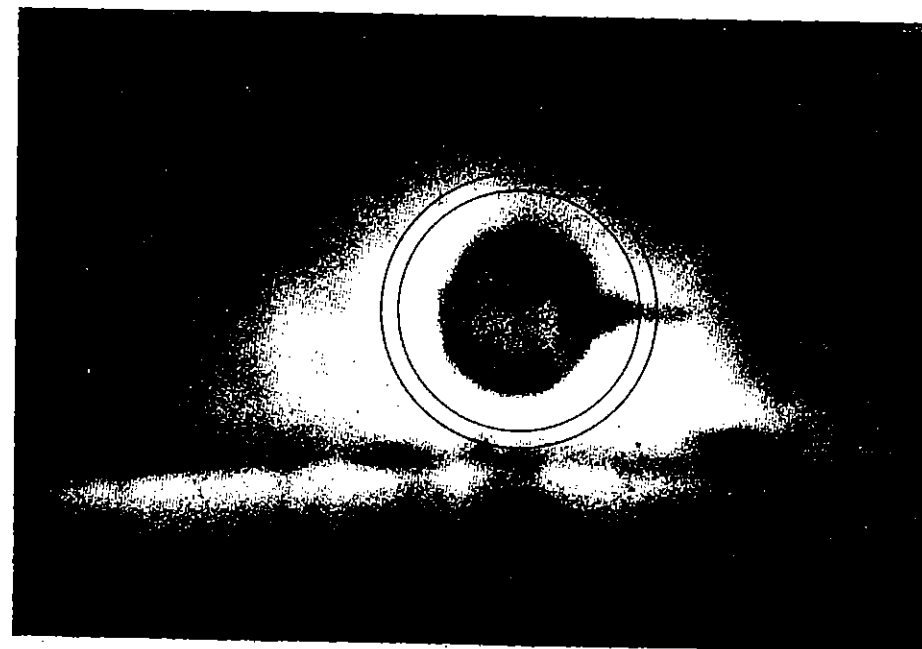


Fig. A2.24 Copper spectrum radiograph of irradiated CH shell.

SHOT 19 ON 19/05/83

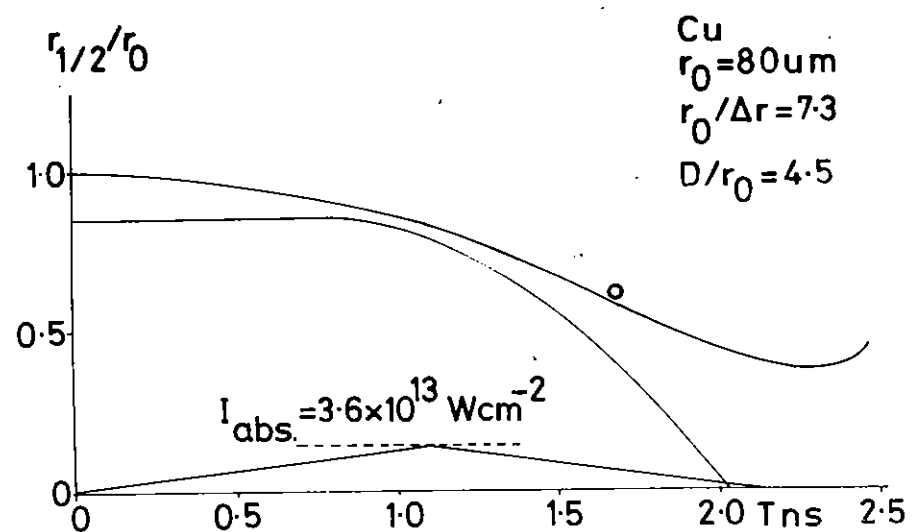
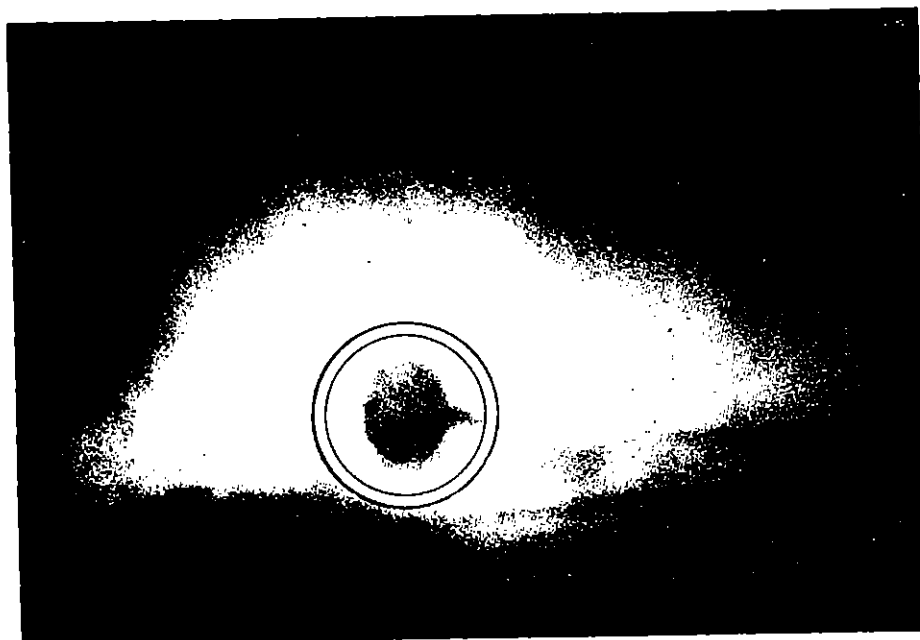


Fig. A2.25 Copper spectrum radiograph of irradiated CH shell.

SHOT 01 ON 20/05/83

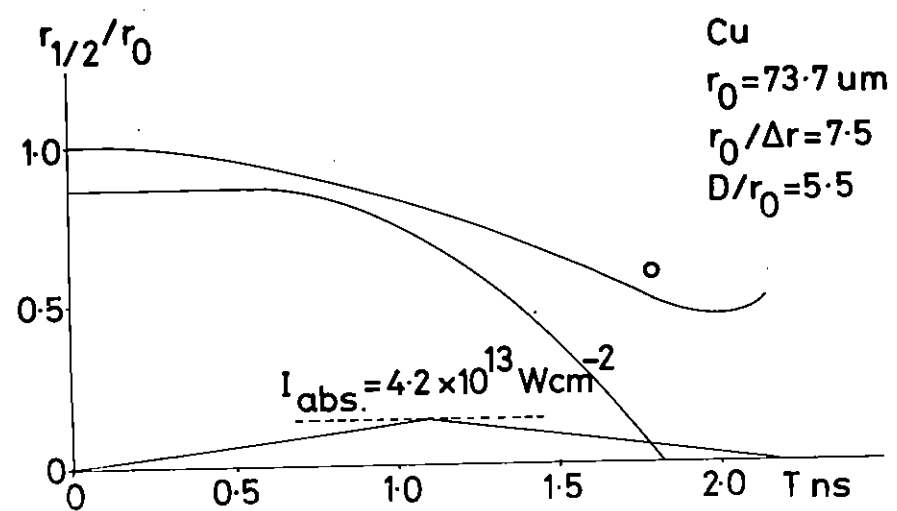
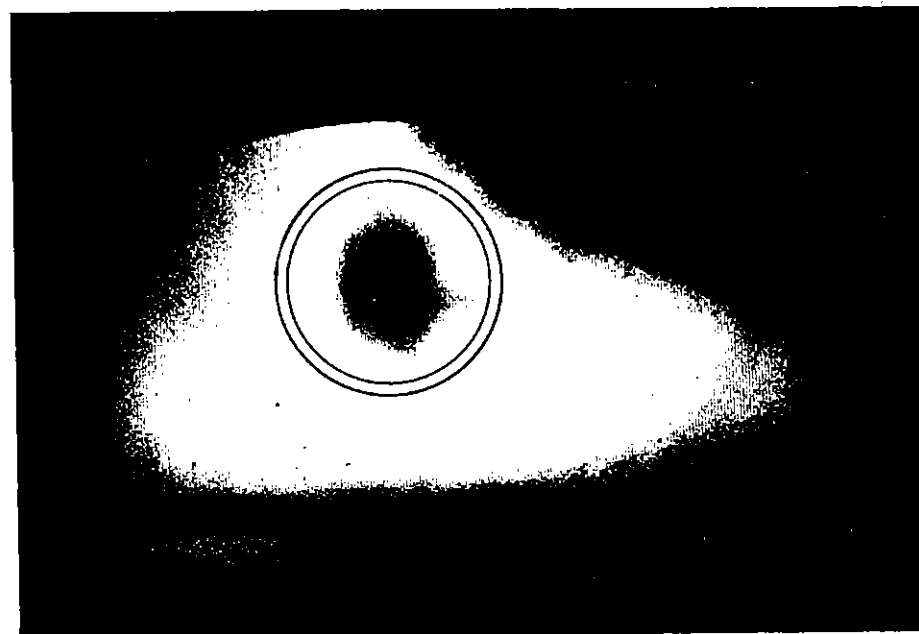


Fig. A2.26 Copper spectrum radiograph of irradiated CH shell.

SHOT 07 ON 08/06/83

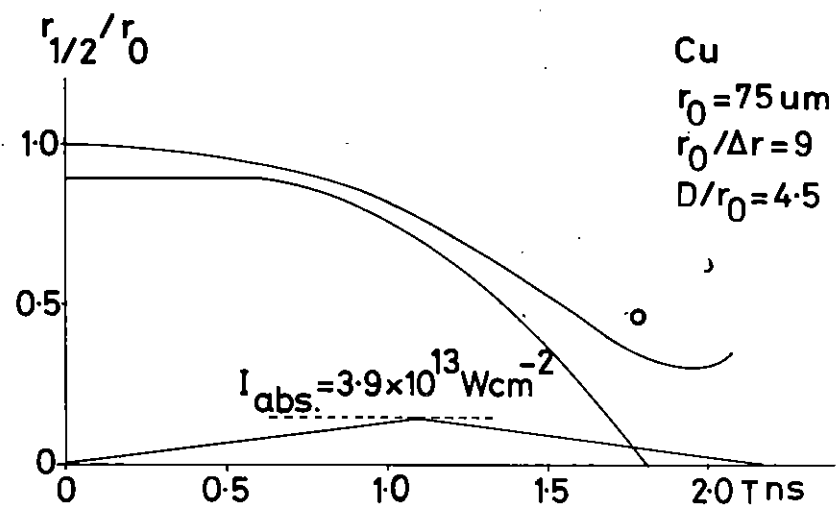
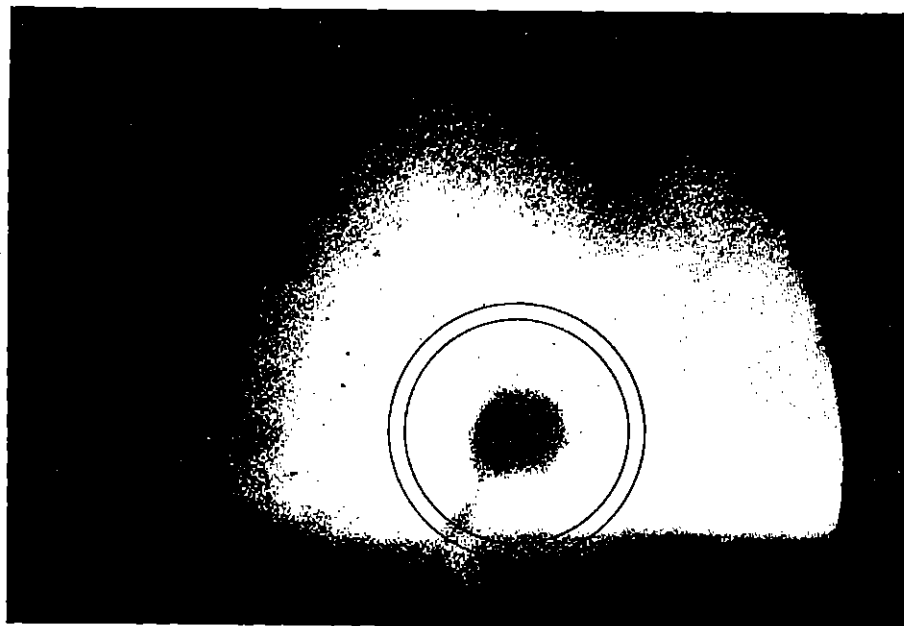


Fig. A2.27 Copper spectrum radiograph of irradiated CH shell.

SHOT 01 ON 10/06/83

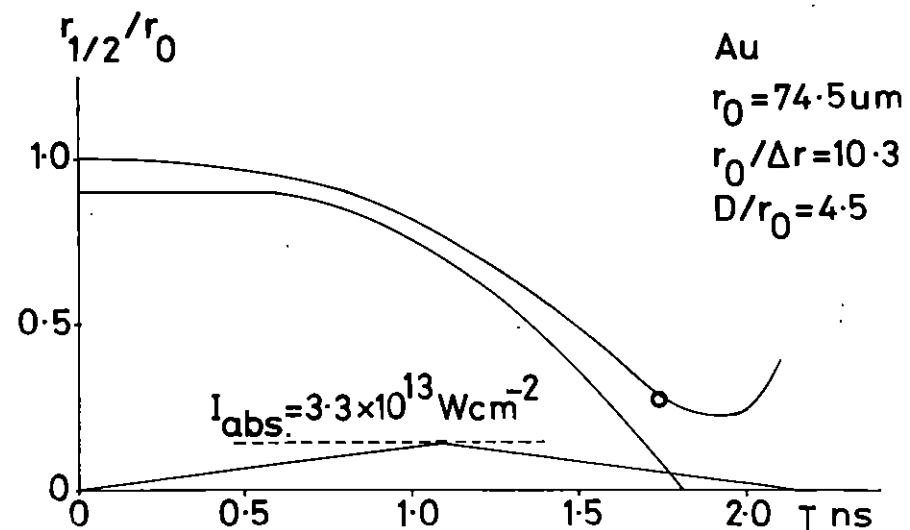
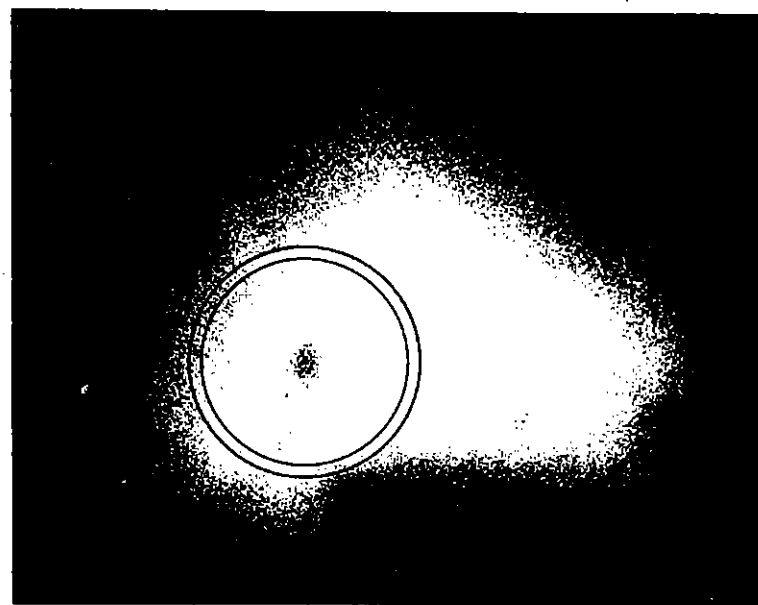


Fig. A2.28 Gold spectrum radiograph of irradiated CH shell.

On one shot, Figure A2.26, where the D/R focussing was changed to 5.5 the hexagonal symmetry appears to have rotated by 30° suggesting a relative enhancement of drive along the double beam overlap directions. (This observation was not reproducible.)

All the Cu backlighting shots show an absorption feature related to the presence of a $\sim 5 \mu\text{m}$ diameter carbon fibre stalk used for mounting the targets. This feature indicates the presence of material, right up to peak compression, spanning the region between the initial stalk contact point and the imploding shell outer surface. The origin of this material is unclear but may be shell material left behind due to the shadowing effect of the plasma produced around the initial stalk which lies in a triple beam overlap direction. Rough estimates of the mass of material required to give the observed opacity and assuming cylindrical symmetry suggest that the material has a density $\leq 1 \text{ g cm}^{-3}$ and comprises a few per cent of the total target mass.

A moderately high volumetric compression with reasonable spherical symmetry has been achieved as seen from the Au radiographs. At peak compression the bulk of the non-ablated target material is confined to a region of radius $(0.2 - 0.25) r_0$ where r_0 is the initial target radius giving a volumetric compression of the order 100.

A2.2.4(e) Data Analysis

The data was quantitatively compared against the predictions of the 1-D hydro-code MEDUSA which was run without radiative preheat. Runs including preheat showed small but negligible effects in the hydro of the shell outer surface for the present ΔR and I_{abs} parameters.

As a comparison standard we choose to define as $r_{1/2}$ the radial distance from the target centre to the chord which would present a line-of-sight attenuation of 50% to the backlighting filtered source spectrum when the intensity was integrated over all frequencies present. The distance $(r_{1/2})_{\text{MED}}$ could be calculated from the appropriate Medusa output time corresponding to the probe time and $(r_{1/2})_{\text{EXP}}$ could be directly measured from

SHOT 03 ON 10/06/83

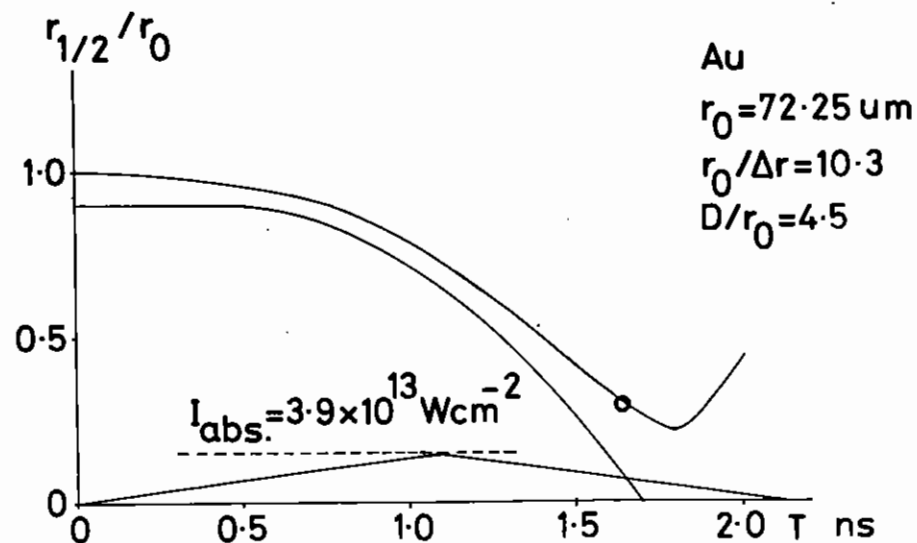
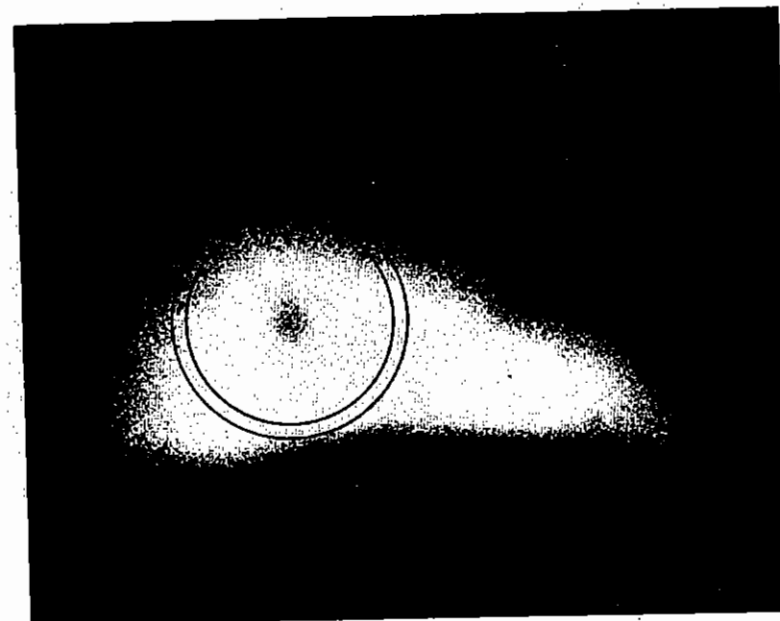
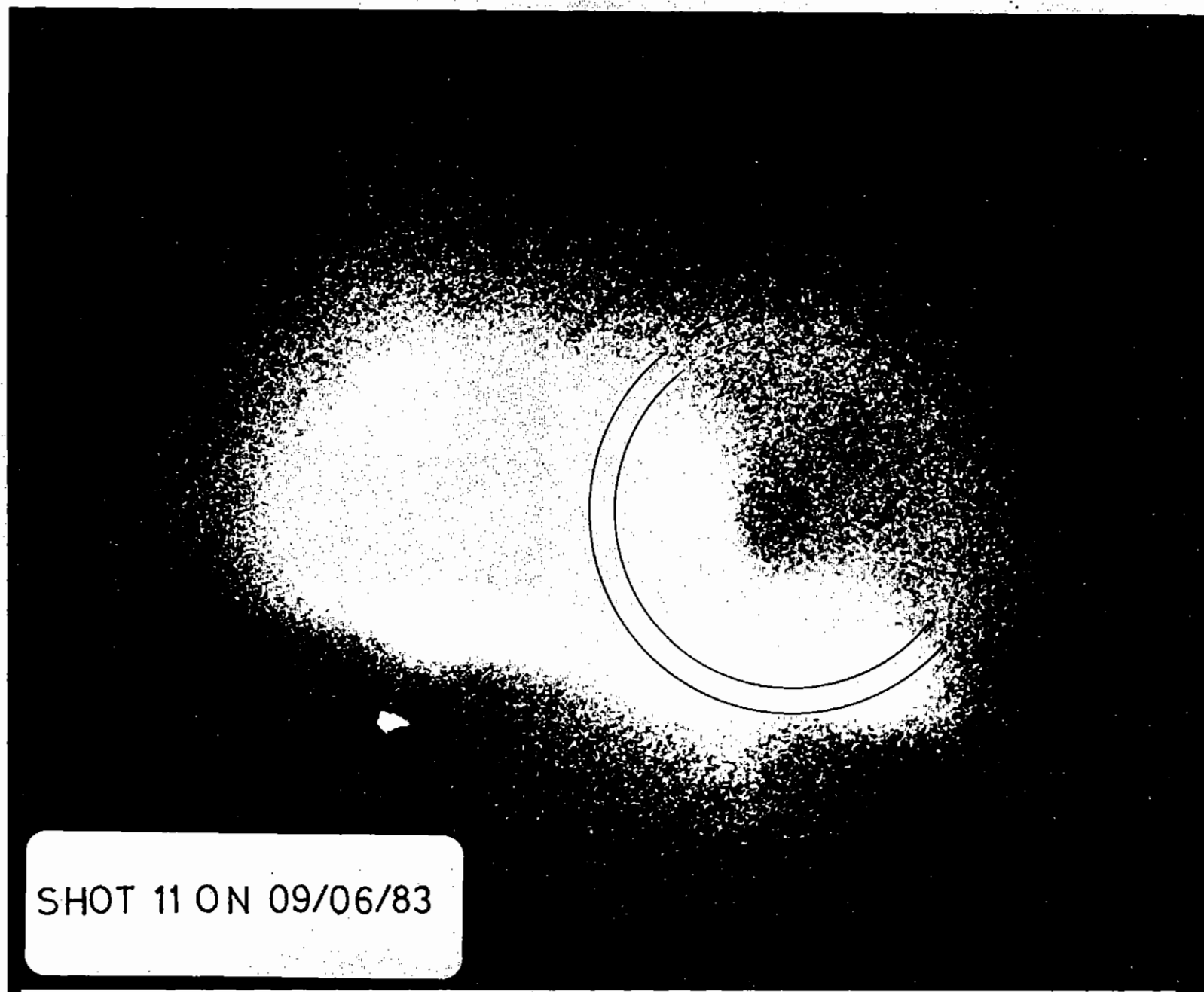


Fig. A2.29 Gold spectrum radiograph of irradiated CH shell.

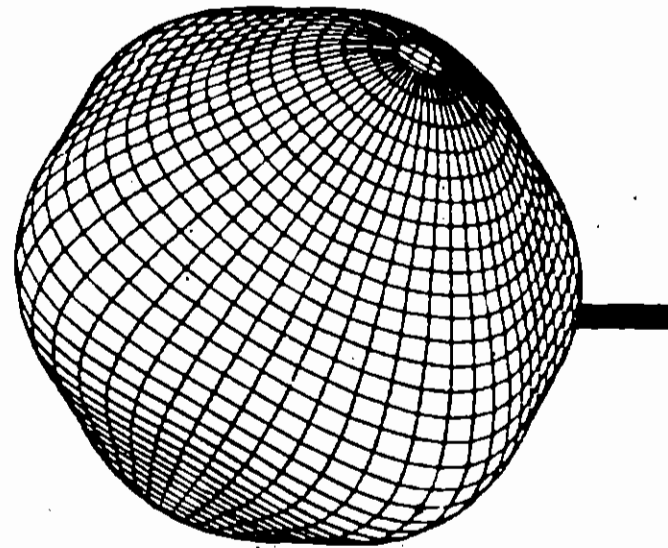


SHOT 11 ON 09/06/83

Fig. A2.30 Gold spectrum radiograph of CH shell target showing shell self-emission.



Shot 19 on 19/05/83



Model

Fig. A2.31 Hexagonal structure as observed in experiment and predicted by modelling.

the raw data using microdensitometry and the detection relative response characteristics.

(r_y)MED was a post-processed quantity from appropriate Medusa outputs giving the material density at each Lagrangian cell position and its evaluation allowed for (a) the polychromaticity of the source spectrum and (b) the convolution of a finite sized imaging pinhole to record the transmitted light.

(i) Source Spectrum Effect

The Medusa output gives the target radial density profile $\rho(r)$ at a given time and a boundary $r = R_0$ beyond which the target material can be taken to be optically thin to X-ray backlighting radiation. The areal mass density along a chord at radius $r = R$ is given by

$$(\int \rho \, dl)_R = \int_R^{R_0} \frac{2r\rho(r)dr}{(r^2 - R^2)^{1/2}}$$

or equivalently by

$$(\int \rho \, dl)_R = 2\rho(R_0)(R_0^2 - R^2)^{1/2} - \int_R^{R_0} 2(r^2 - R^2)^{1/2} (d\rho(r)/dr)dr$$

where the first term is finite only if $\rho(R_0) > 0$ as in the case of a cold target with a discontinuity in its density profile.

Hence the transmission profile of the target to monochromatic backlighting is given by

$$T_E(R) = \exp(-2\mu_E \int_R^{R_0} (r^2 - R^2)^{1/2} (d\rho(r)/dr) dr)$$

when $\rho(r > R_0) = 0$ and μ_E is the target material mass attenuation coefficient at the probe radiation energy, E .

For a broad-band source spectrum the transmission profile is

$$T(R) = \int_0^\infty I_E T_E(R) dE / \int_0^\infty I_E dE; R < R_0$$

$$= 1.0; R > R_0$$

where I_E is the relative intensity of the source spectrum after allowing for fixed filters in the detection channel. Transmission profiles were calculated from Medusa outputs for CH targets assuming cold material absorption coefficients and using $\mu_E = 2000/E^3 \text{ cm}^2 \text{ g}^{-1}$ above the carbon K-edge with the photon energy E in KeV.

(ii) Pinhole convolution effect

The function $T(R)$ becomes the relative radial source brightness function in the target plane for projection via a pinhole of diameter D onto the detection plane. The relative transmission function as recorded in the final image plane is calculated as follows with reference to Figure A2.32. A point at radius r' in the detector plane will receive intensity from all points in the source (shell target) plane within the shaded circle of area $\pi D^2/4 (1 + 1/m)^2$ where $m = v/u$ is the pinhole camera magnification. All relevant ray bundles are sufficiently close to the principal axis that angle projections are ignored. Radial variations in source brightness throughout the region of the shaded circle is allowed for by calculating the relative transmission profile in the detector plane as

$$T(r') = \int_a^b T(r) 2r\theta(r) dr / (\pi D^2/4 (1 + 1/m)^2)$$

where

$$a = r'/m - D/2 (1 + 1/m)$$

$$b = r'/m + D/2 (1 + 1/m)$$

$$\theta(r) = \cos^{-1} \left(\frac{(r'/m)^2 + r^2 - (D/2 (1 + 1/m))^2}{2rr'/m} \right)$$

and valid for $r' > D/2 (m + 1)$.

For $r' = 0$ the transmission is given simply by

$$T(0) = \int_0^{D/2(1 + 1/m)} 2\pi r T(r) dr / (\pi D^2/4 (1 + 1/m)^2)$$

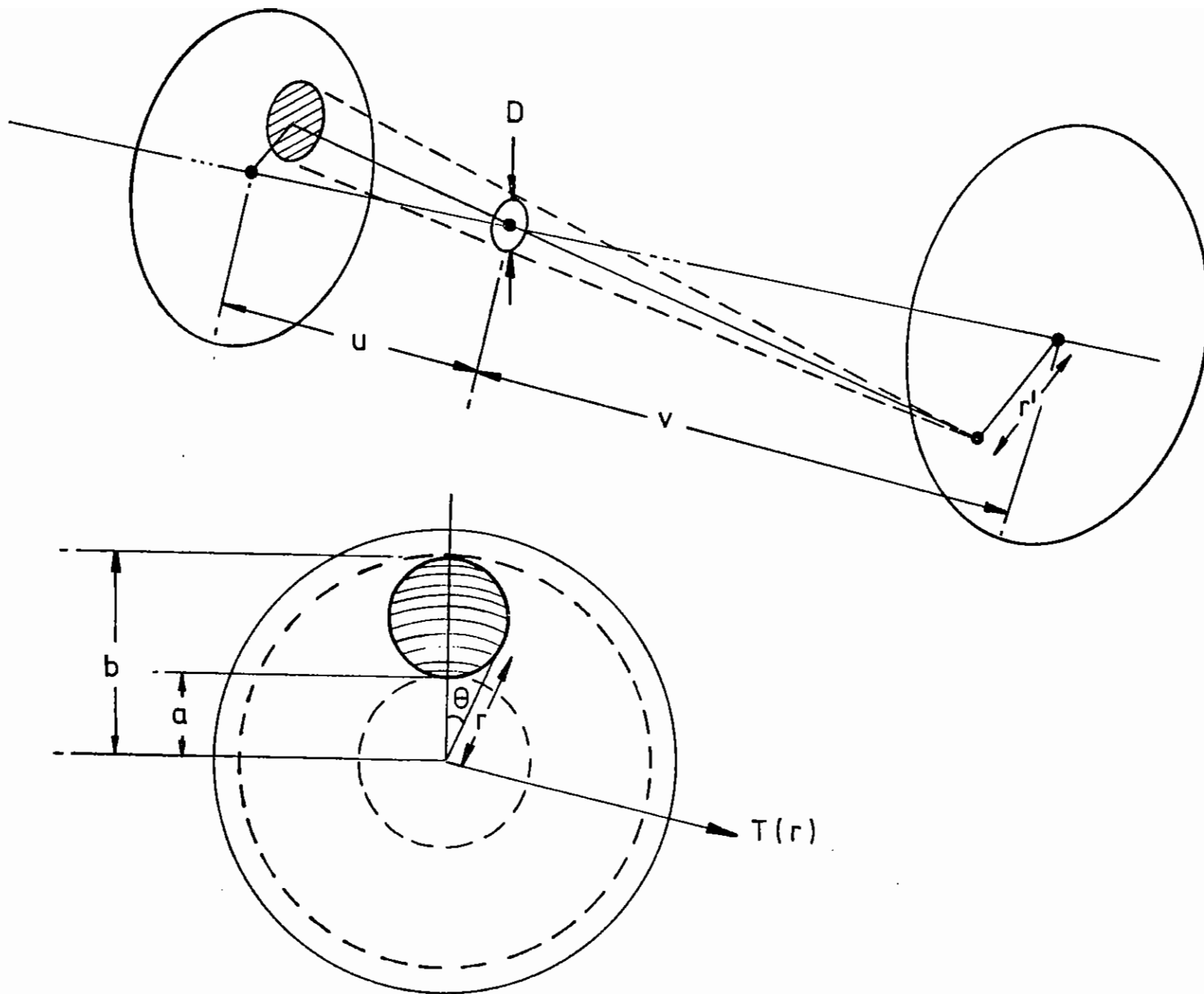


Fig. A2.32 Geometry of pinhole convolution procedure.

This profile can be directly compared to transmission profiles deduced from raw data.

Figure A2.33 shows the measured and calculated transmission profiles corresponding to the cold shell radiograph of Figure A2.21. This fit gives the effective spatial resolution $D \sim 15 \mu\text{m}$, the overall magnification $m = 50$ and indicates the validity of the backlighting spectrum used in further analysis. Figure A2.34 shows the measured and calculated transmission profiles for an Au radiograph near peak implosion (cf Figure A2.28) using the pinhole diameter $D = 20 \mu\text{m}$.

Figures A2.22 to A2.29 also show the Medusa predicted trajectories of the $r_{1/2}$ chord during the implosion (as well as the shell inner boundary) and also the experimental observation at the probe time. In all cases the agreement is good and is summarised in Figure A2.35 where the normalised $r_{1/2}$ ratios are plotted against a normalised time parameter t_p/t_{imp} and where t_p is the probe time and t_{imp} is the computed peak implosion time for a given target. The Cu data close to peak implosion time is less reliable due to the longer frame time of this soft energy channel and also due to the more sensitive nature of the transmission profile to softer X-rays in the outer plasma when cold material absorption coefficients are probably inadequate.

The good agreement between Medusa and experimentally measured hydro movement gives indirect evidence that the targets achieved final core parameters comparable to theoretical predictions. In the case of the Au radiographs where a complete radial transmission profile was measurable a secondary check was possible by carrying out a deconvolution of the pinhole diameter D and an Abel inversion procedure to obtain a radial density profile of the core. A rough deconvolution procedure was followed by transforming the measured transmission profile, $T(r)$, to a deconvoluted transmission profile $T'(r)$ with the following conditions satisfied.

$$1) \quad T(r_{1/2}) = T'(r_{1/2}) = 0.5$$

$$2) \quad T(r_0) = T'((r_0 - (D/2)^2)^{1/2}) = 1.0$$

$$\text{and } 3) \quad \int_0^\infty (r(1 - T(r))dr = \int_0^\infty r(1 - T'(r)) dr$$

An Abel inversion, allowing for the spectral content of the filtered backlighting source was then carried out on the function $T'(r)$ by the method described in the previous annual report (A2.17). Figure A2.36 shows a deduced density profile near peak compression from the raw data of Figure A2.28 together with Medusa density profiles at each end of the radiograph frame time over which the data will be averaged. Again the agreement is good indicating that peak densities of the order $\rho = 10 \text{ g cm}^{-3}$ have been achieved. The hydrodynamic efficiency of these implosions was estimated to be typically $\sim 10\%$ using a maximum shell velocity of $\sim 5 \times 10^6 \text{ cm/sec}$ from medusa shell trajectories.

A2.2.4(f) Conclusions

A set of target and illumination parameters have been established for the currently available output of the VULCAN laser system which give rise to implosion dynamics consistent with modelling by the simulation code Medusa. For experimental comparisons simulations did not attempt to allow for effects such as radiative preheat, Rayleigh-Taylor instability or illumination non-uniformity indicating that for the present experimental conditions the presence of these effects does not seriously degrade the quality of the implosion. The major factor in achieving these good implosions and high final density has been the choice of low aspect ratio targets. This is illustrated in Table A2.3 and Figure A2.37 where the current experimental parameters are put in the context of similar experiments carried out in recent years and reported in previous annual reports. It is clear there is a strong correlation between aspect ratio, agreement with simulation and peak density attained.

Although this work has not been quantitative in its assessment of the adverse effects expected to be important in achieving higher densities it is anticipated that future work with twelve beams will allow higher aspect ratio targets driven with higher, more uniform ablation pressure to reach

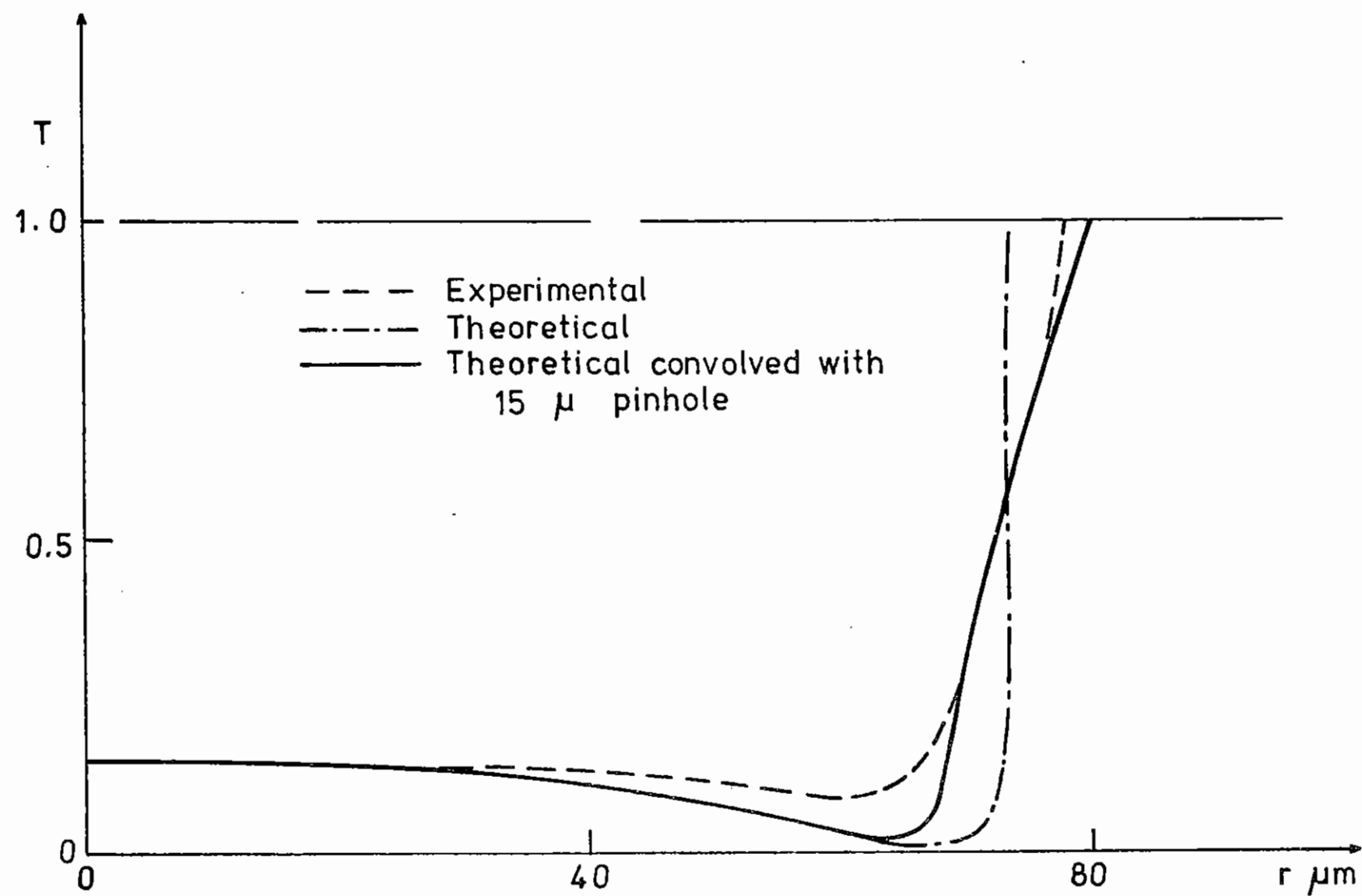
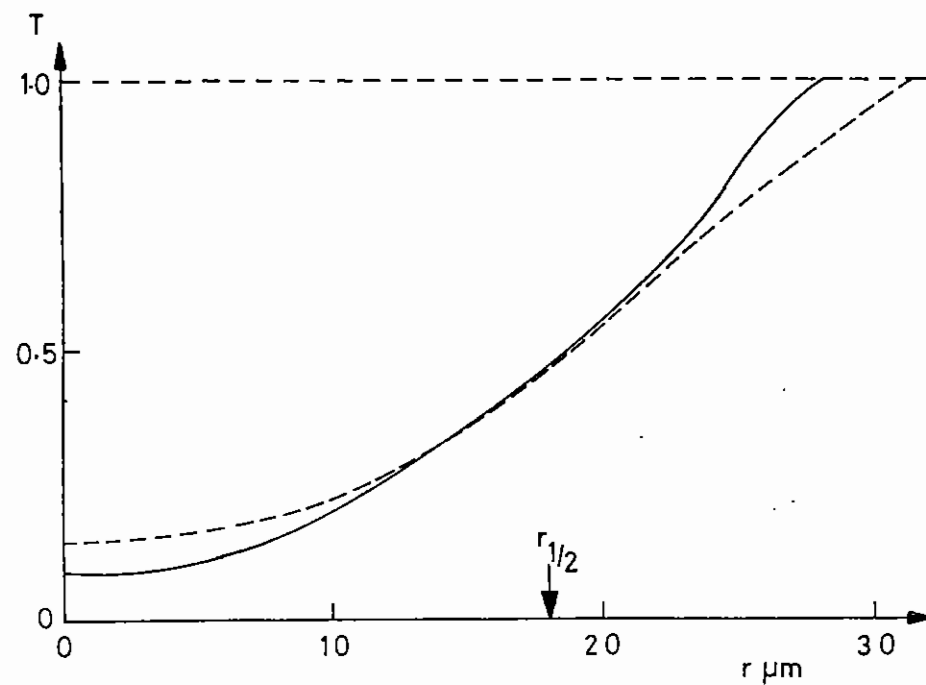
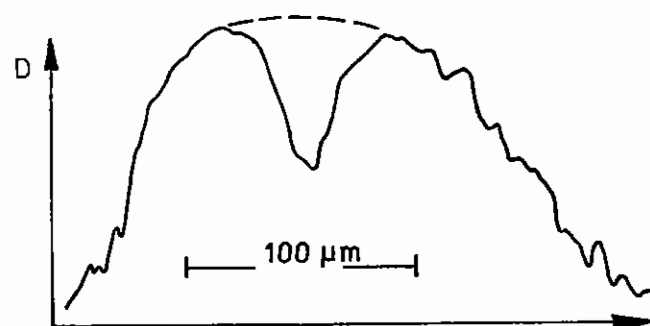


Fig. A2.33 Fitting of theory to experiment for a cold shell radiograph.

SHOT 01 ON 10/6/83



— Experimental Transmission
 --- Transmission convolved from
 Medusa (20 μm P/h)

Fig. A2.34 Theory and experimental agreement for a gold radiograph near peak compression.

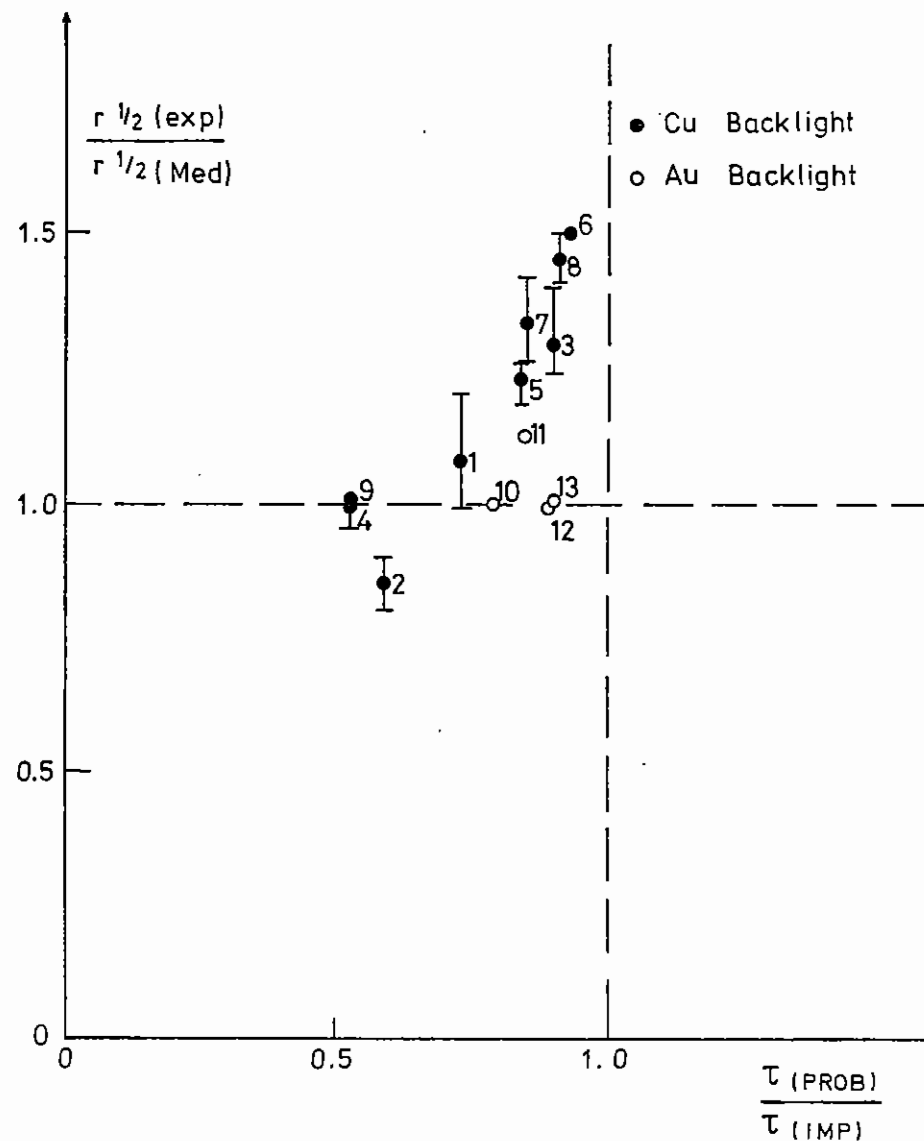


Fig. A2.35 Data comparison to Medusa calculations. Symbols on points refer to Table A2.2.

Shot 01 10 / 6 / 83

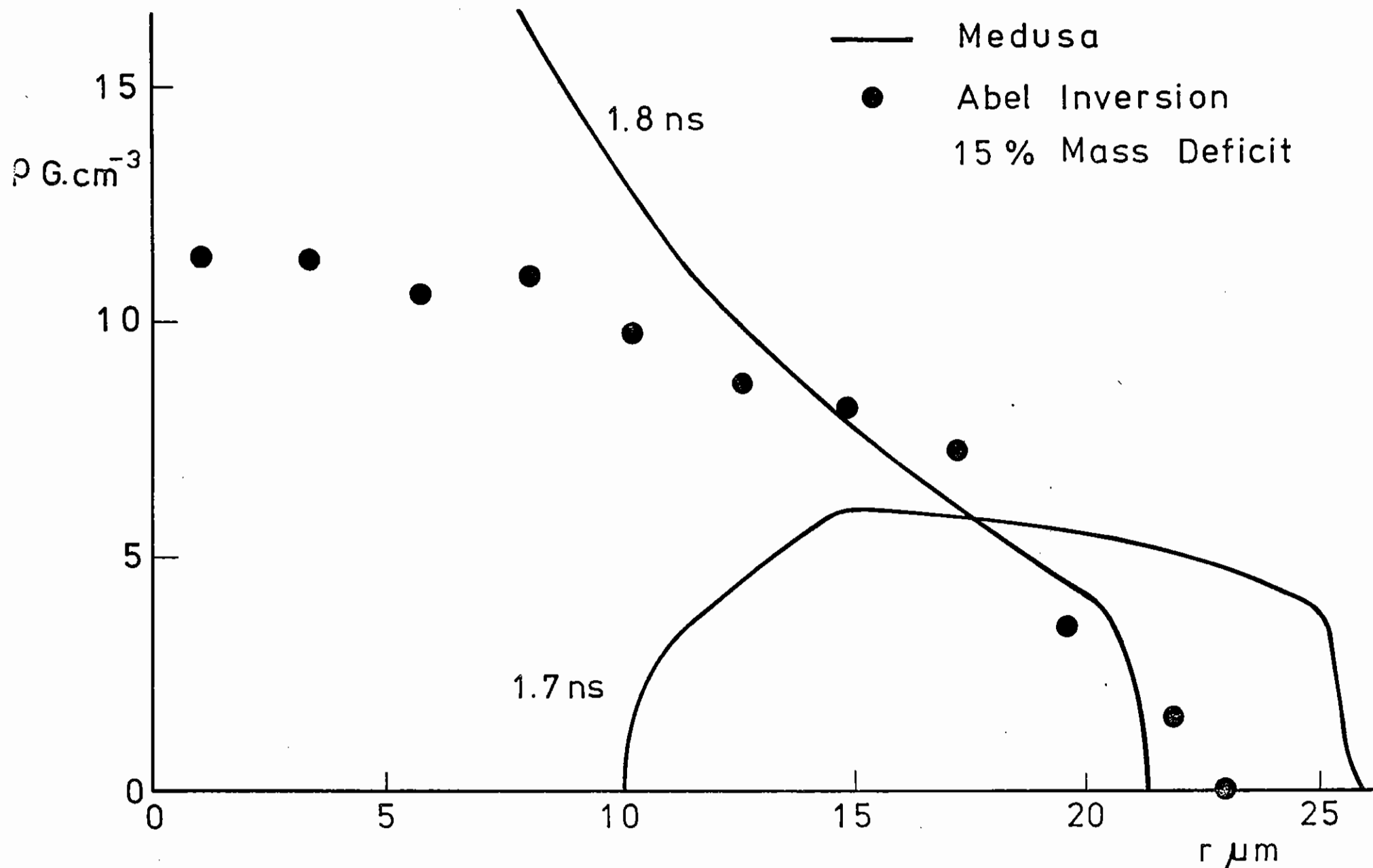


Fig. A2.36 Density profile deduced from Abel inversion of a gold radiograph.

λ μm	0.53 ^o	0.53 [•]	0.53 ^Δ	0.53 [□]
I_{ABS} 10^{13} W cm^{-2}	2-4	1.5-3	0.5-2.5	≈ 1.0
τ_L ns	0.9	1.3	1.0	0.5
$r/\Delta r$	7-10	5-7	20-60	≈ 20
r μm	75-90	60-75	75-95	≈ 85
D/r	4.5	3.8	4.2	3.6
E_{probe} keV	1.6/2.5	2.0	1.7	1.5
Method	PULSED	PULSED	STREAKED	PULSED
ρ_{exp} g cm^{-3}	≈ 10	≈ 5	≈ 1	≈ 1
SYMBOL REFERENCE FOR FIGURE A2.37	o	•	Δ	□

TABLE A2.3

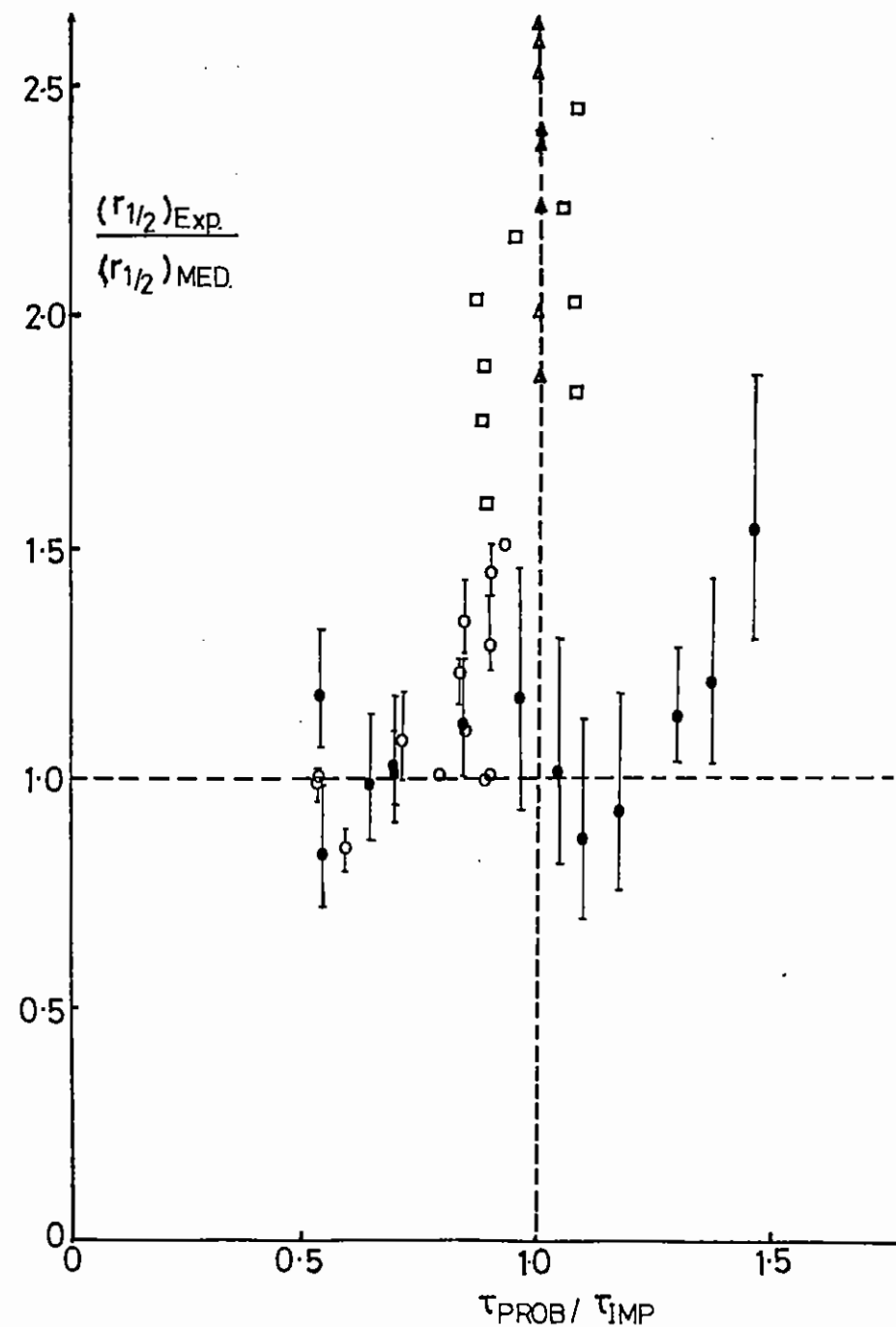


Fig. A2.37 Present and earlier data comparison to Medusa calculations. Symbols refer to Table A2.3.

higher peak density. To take full advantage of more uniform illumination it may be necessary to eliminate mounting stalk perturbations by using unsupported targets and to monitor higher ρv values by using higher energy backlighting sources. Preliminary investigations into inertially suspended targets and high energy, spectrally resolved backlighting techniques have already begun.

A2.2.5 Emission Diagnosis of 0.53 μ m Ablative Implosions

B J MacGowan, J D Kilkenny (IC), S J Rose and M H Key (RAL)

Initial experiment design revolved around the utilization of the polymer microballoon targets filled with Freon 12 to produce hot emitting cores surrounded by the remains of a low opacity polymer shell (A2.18). Low aspect ratio target designs featuring the polymer shells have been described in section A2.2.2. Due to poor energy coupling into the targets as a result of the large amount of defocus required for uniformity, the absorbed energy was insufficient to drive an implosion fast enough to produce a hot emitting core in a low aspect ratio target.

Space resolved spectroscopy was achieved using a standard minispectrometer with a 10 μ m slit positioned so as to give $\times 6$ magnification. The resulting spectra taken on SB5 film showed no core emission from the thicker walled polymer balloons. The best result was achieved when the aspect ratio constraints were relaxed and a balloon of wall thickness 4.2 μ m and aspect ratio 16 was imploded. The spectrum in the region of the Cl XVI He α is shown in Figure A2.38 together with the associated pinhole image. The Li-like satellites to the He α are just resolvable and indicate a density peak in the centre. The high background noise level in this spectrum makes more quantitative analysis difficult.

Further work was done on implosions of D-T filled glass balloons. These targets were unsatisfactory from the designers point of view due to the large aspect ratios (as much as 20). Typical diameters were 150-160 μ m and wall thicknesses 1.5-3 μ m. Neutron yields of up to 10^6 were observed and when compared with a constant temperature, constant density model implied

core ion temperatures in the range 200 - 500eV. Figure A2.39 shows the inferred ion temperatures plotted against the calculated peak ρv^2 of the target wall. The ρv^2 was calculated using the absorbed irradiance and ablation pressure scaling laws derived from experiment at RAL (A2.19). There seems to be some correlation with ρv^2 . Points are also shown which were obtained in December 1981 at RAL (A2.8). The targets used in December 1981 were slightly thinner walled but do not deviate significantly from the more recent data.

Space resolved spectra were obtained from the D-T filled glass microballoons. An example of the silicon spectra obtained is shown in Figure A2.40. The presence of a Calcium impurity in the glass was detected in 2nd order Bragg diffraction, from the core region of the target. Emission to the long wavelength side of the Ca He α may be indicative of less ionised states of Calcium emitting 2 - 1 (K α) radiation.

A2.2.6 Thermonuclear Spectroscopy with CR39 Track Detectors

A P Fews, D L Henshaw (Bristol), D A Pepler, S J Rose and W T Toner (RAL)

Detailed analysis of α -particle and proton tracks in CR39 prepared and etched under carefully controlled conditions enables the energy of the particles to be determined with a precision approaching the straggling limit (A2.20). This technique has been applied to ablatively driven implosions of DT filled glass-walled microballoons. With thin walled balloons ($t < 1.6 \mu$ m) both α -particles and protons penetrated the stagnated shell material to be registered in this detector. Proton tracks have been recorded for implosions with balloon wall thickness of up to 3.2 μ m.

CR39 detectors were positioned 30 mm from the target, arranged to occupy the maximum possible available solid angle. Mylar layers of various thickness from 2 to 60 μ m were used to protect the CR39 on several of the shots. The microballoons had a gas fill of 20 bars DT. The beams were focused using a d/R parameter of 4.5.

The first step in the determination of the range, and hence the energy, of

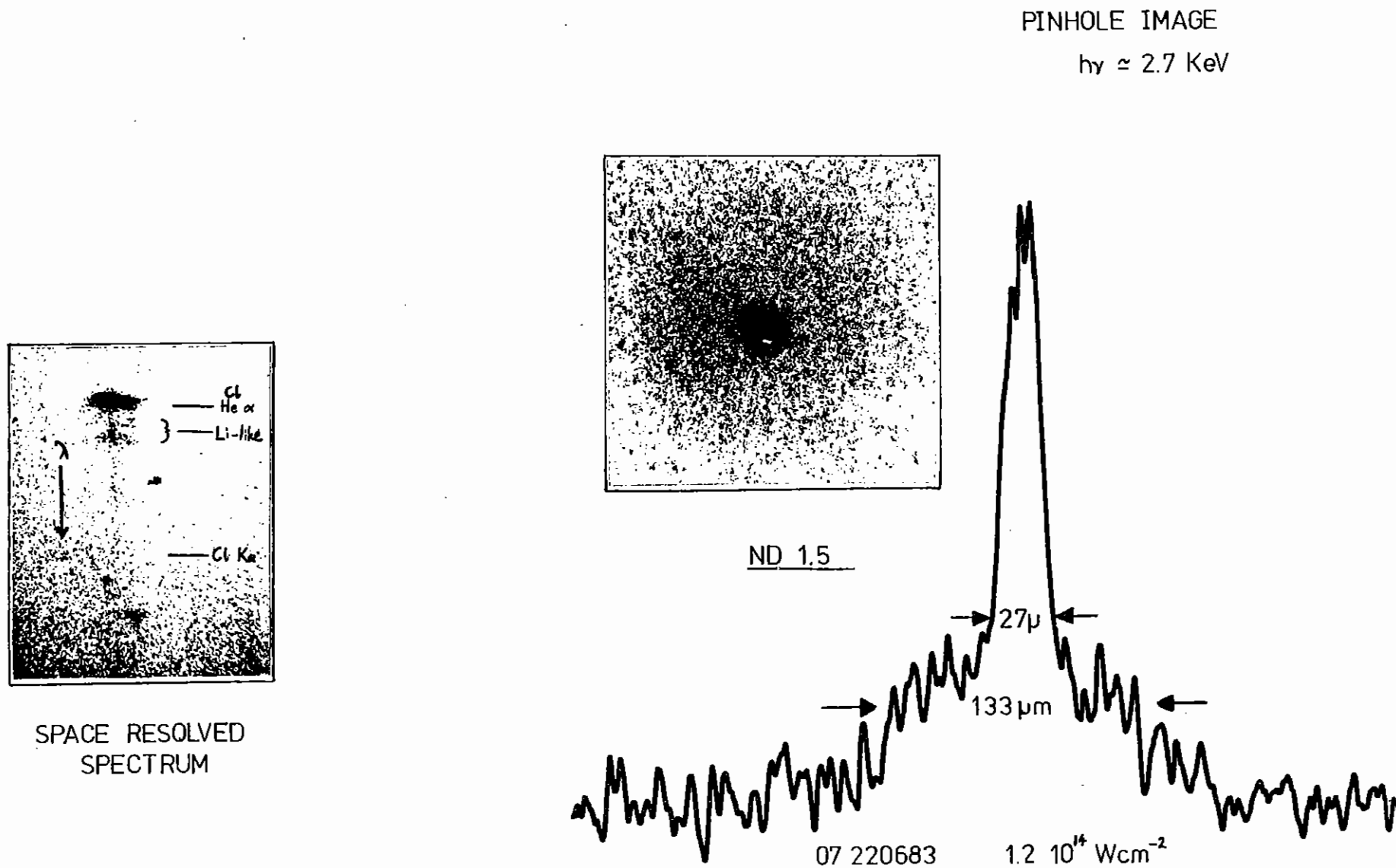
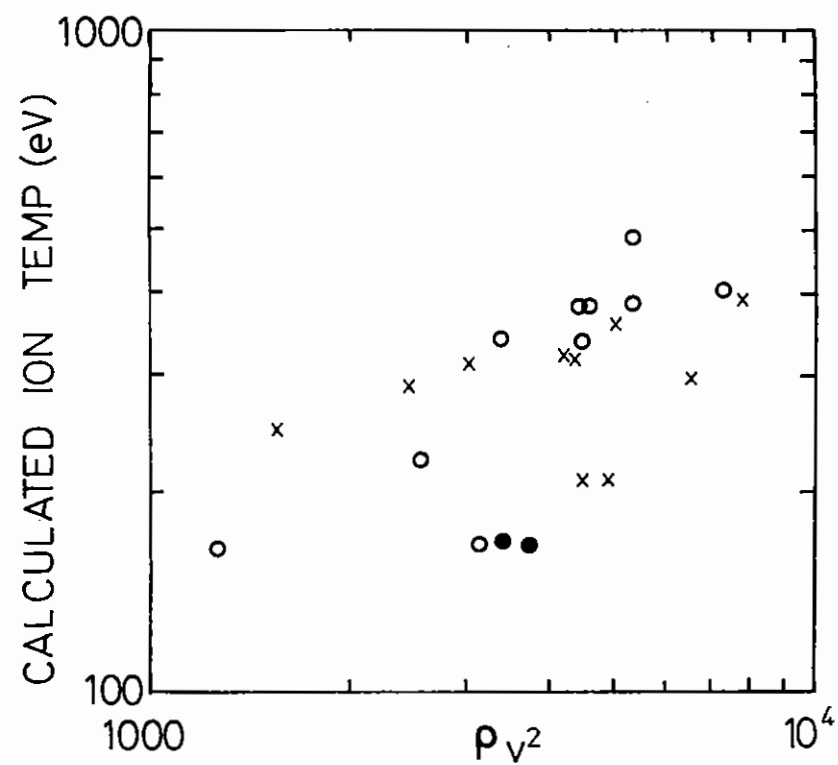


Fig. A2.38 Space resolved spectrum and pinhole picture for high aspect ratio polymer microballoon filled with 1 bar Freon 12. The central target diameter was 133 μm and wall thickness 4.2 μm .



DARK CIRCLES JUNE 83 Δr_g 1.6—3.3 μm

'x' DEC. 81 Δr_g 1.0—2.0 μm

Fig. A2.39 Estimated ion temperature based on measured neutron yield for compression carried out in December 1981 and June 1983.

SHOT 19 230683

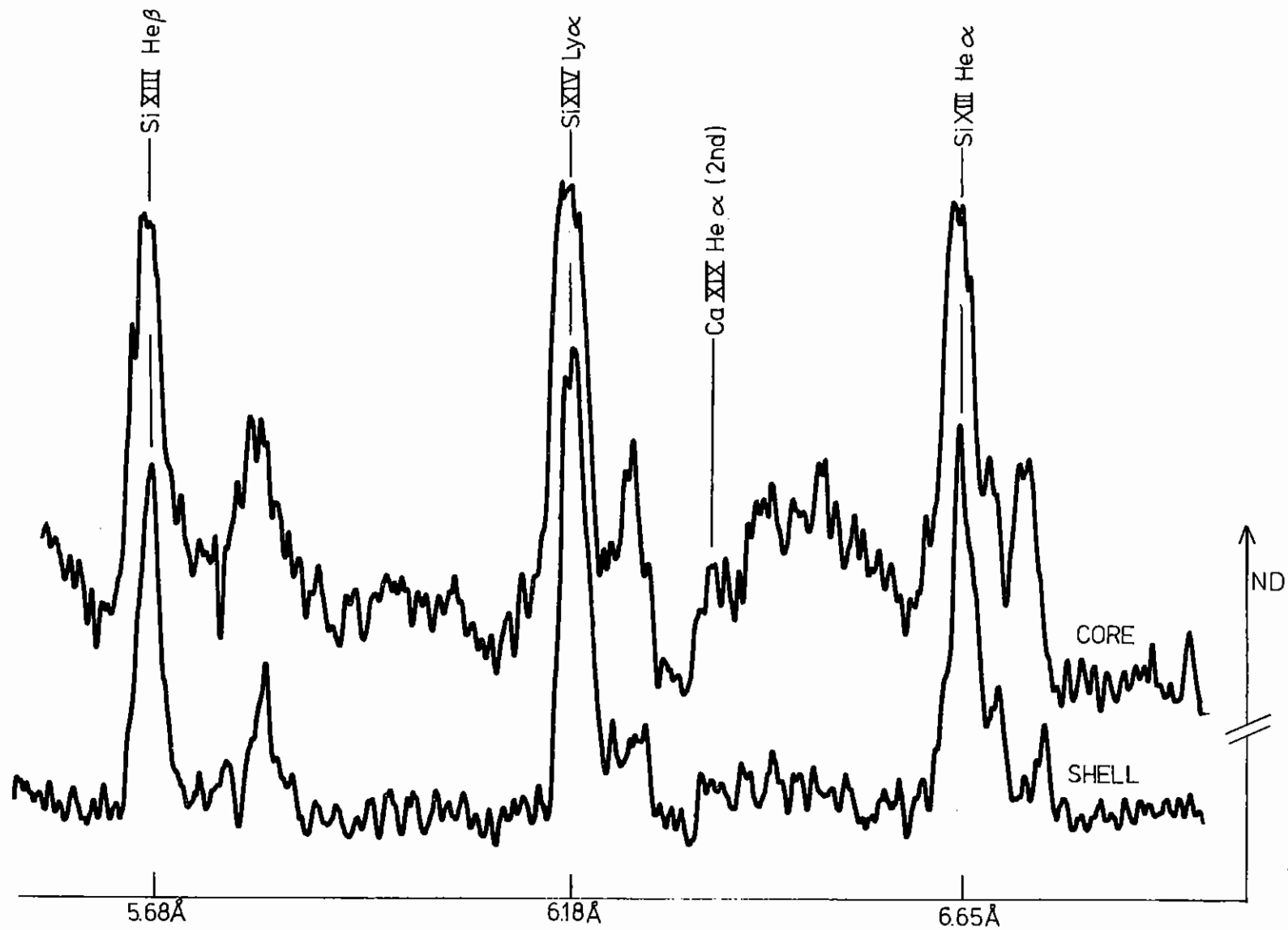


Fig. A2.40 Space resolved spectrum of a glass microballoon. This shot at absorbed irradiance $10^{14} \text{ W cm}^{-2}$ produced $7 \cdot 10^5$ neutrons. The balloon was 158 μm in diameter and 1.6 μm wall thickness.

the incident particle is to measure the dimensions of the etch pit under the optical microscope. The best estimate of the particle range computed from these measured parameters includes compensation for range straggling and multiple coulomb scattering within the CR39 itself. This is described in more detail in A6.7.2.

The track response of the CR39 is calibrated by measuring the ratio of the track etch rate ($\frac{V_T}{V_B} - 1$) as a function of particle range for both α -particles and protons. The track etch rate rises with decreasing particle range, reaching a Bragg type peak at $\sim 1.9 \mu\text{m}$ (0.37 MeV) for α -particles and $\sim 0.7 \mu\text{m}$ (0.07 MeV) for protons.

Figure A2.42 shows a photomicrograph of α -particle tracks and a proton track from shot 14, 23/6/83 given in the table of results (Table A2.4). The proton track appears substantially darker than the α -tracks because it has a range of $40 \mu\text{m}$ and consequently still has a cone-like structure. The α -tracks are heavily over-etched and only a spherical remnant of each track remains.

Figure A2.43 shows examples of the measured α -particle spectra from shots 14 and 19 on 23/6/83, with glass wall thicknesses of $1.2 \mu\text{m}$ and $1.6 \mu\text{m}$ respectively. No mylar protection was used for either shot. The resolution of the CR39 for these examples is estimated to be $\sim 0.1 \mu\text{m}$.

Figure A2.44 shows two examples of measured proton spectra, from shot 19, 23/6/83 which may be compared with the corresponding α -particle spectrum in the previous figure and for shot 18, 21/6/83. The number of tracks reaching the CR39 detectors are limited because the (DD) proton yield is only 1% of the DT reaction yield. The difference in α -particle and proton density radius product (ρR) for shot 19, 23/6/83 is due to systematic error in the proton response calibration.

Table A2.4 shows a summary of the principal results obtained. The neutron yield was monitored with a plastic scintillation counter. The neutrons and α -particles, both produced in the DT reaction are present in equal abundance

and the protons are produced by the DD reaction at typically 1% of the DT reaction yield. Both the initial balloon wall thickness and the thickness measured from the thermonuclear products (ρR) are given in units of $10^{-3} \text{ g cm}^{-2}$. The $\sigma(\rho R)$ gives the standard deviation in the measured ρR . Ion temperatures deduced from Neutron yields are shown in Figure A2.45 as a function of specific absorbed energy. A uniform temperature and density was assumed.

A graph showing the variation of ρR with initial microballoon wall thickness is shown in Figure A2.46. It can be seen that the target thickness ρR measured at the time of thermonuclear yield is between 5 and 11 times the initial wall thickness.

Discussion

Since the conditions of the radiographic shots could not be reproduced exactly in this phase of the experiment, a set of targets of progressively increasing wall thickness was chosen so that trends seen in the data could be extrapolated for comparison. The ion temperatures deduced from the yields shown in Figure A2.45 are in good agreement - perhaps fortuitously so - with the results of two simulated shots of thick walled targets at a rather higher irradiance than was actually used. But the data show a smooth progression indicating that all targets behaved in a similar manner. The uncertainties ($\sim \pm 25\%$) in relative DD/DT yield mean that the yield ratio can only serve as a consistency check on the temperature. The observed ratios of a factor of 100 on average suggest a slightly higher temperature ($\sim 500 \text{ eV}$) than either the absolute yields or the widths of the alpha and proton ρAR distributions would suggest. The mean values of the ρAR distributions for alphas and protons show that thermonuclear emission takes place when the shell ρAR has increased between four-fold and eleven fold which is substantially less than the radiographic data give at peak compression. The most remarkable feature of the data is the small values observed for the widths of the α and proton ρAR distributions which are never more than a few percent of the mean ρAR . The following effects all contribute to $\sigma(\rho R)$:

Shot Index	Absorbed Energy (J)	4 π neutron yield	T _i (n=100) eV	4 π α yield	4 π proton yield	Number of protons measured	Balloon diameter (μ m)	Wall thickness μ m	10^{-3} .g.cm ⁻²	α -particle $\rho\Delta R$ 10^{-3} .g.cm ⁻²	α -particle $\sigma(\rho\Delta R)$ 10^{-3} .g.cm ⁻²	proton $\rho\Delta R$ 10^{-3} .g.cm ⁻²	proton $\sigma(\rho\Delta R)$ 10^{-3} .g.cm ⁻²
14,23/6/83	34	10 ⁶	403	1.3 $\pm$ 0.1 $\times$ 10 ⁶	N/A	3	155	1.2	0.28	1.616 $\pm$ 0.005	0.061	2.2 $\pm$ 0.1	0.24
19,23/6/83	37	6.5 $\times$ 10 ⁵	377	4.7 $\pm$ 0.5 $\times$ 10 ⁵	5.6 $\pm$ 1.4 $\times$ 10 ³	20	158	1.6	0.37	1.770 $\pm$ 0.003	0.048	2.3 $\pm$ 0.1	0.18
11,24/6/83	35	1.8 $\times$ 10 ⁵	337	-	7.0 $\pm$ 2.8 $\times$ 10 ²	4	142	1.9	0.43	-	-	4.0 $\pm$ 0.2	0.38
15,21/6/83	50	7.4 $\times$ 10 ⁴	303	-	3.4 $\pm$ 1.7 $\times$ 10 ²	3	140	2.5	0.58	-	-	5.1 $\pm$ 0.3	0.36
18,21/6/83	45	6.7 $\times$ 10 ⁴	280	-	5.8 $\pm$ 1.7 $\times$ 10 ²	12	161	2.7	0.62	-	-	5.5 $\pm$ 0.1	0.25
23,23/6/83	34	4 $\times$ 10 ³	170	-	2.7 $\pm$ 2.7 $\times$ 10 ¹	1	163	2.9	0.67	-	-	7.4	N/A
19, 5/7/83	22	7 $\times$ 10 ³	210	-	8.1 $\pm$ 4.6 $\times$ 10 ¹	3	126	3.2	0.74	-	-	8.3 $\pm$ 0.4	0.73

Table A2.4

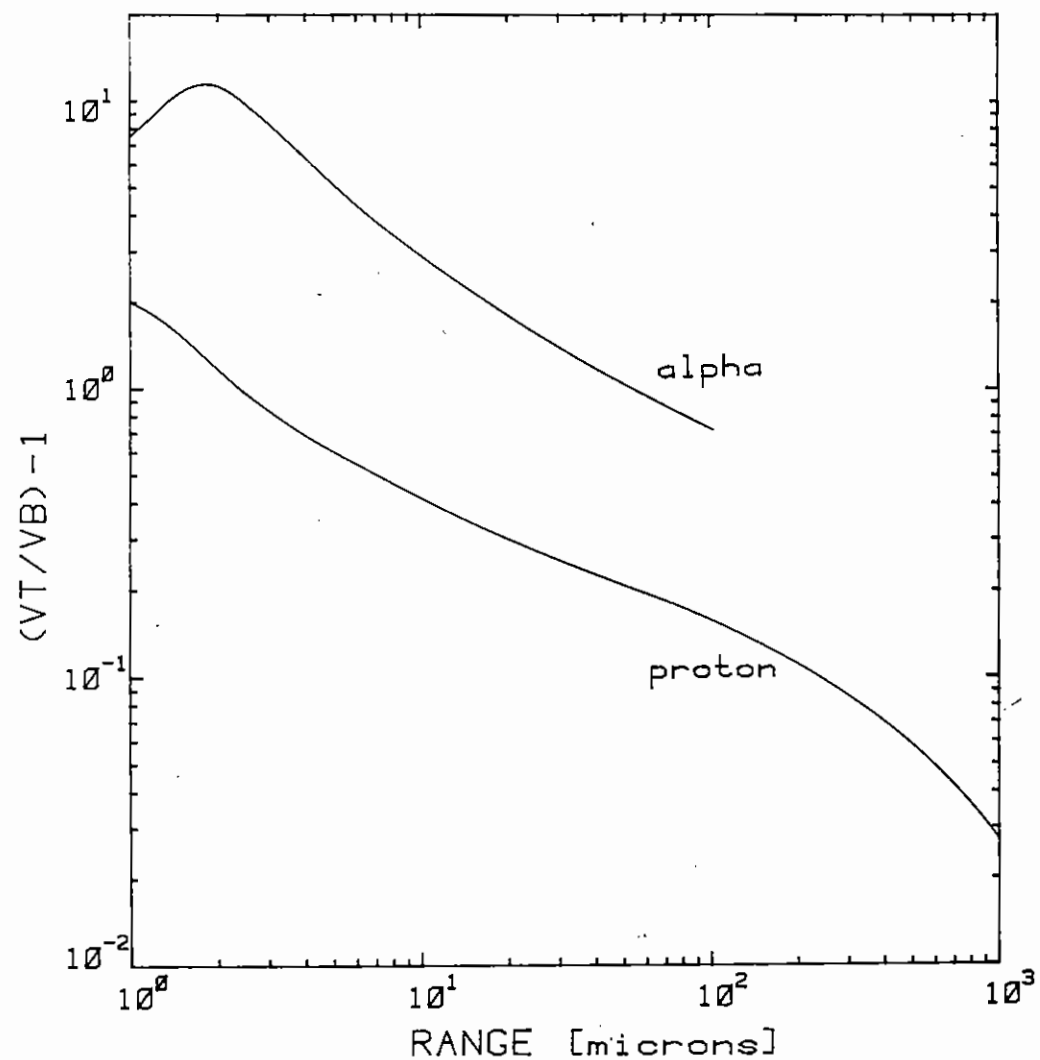


Fig. A2.41 The variation of etch sensitivity $(V_T/V_B - 1)$ versus particle range for α -particles and protons.

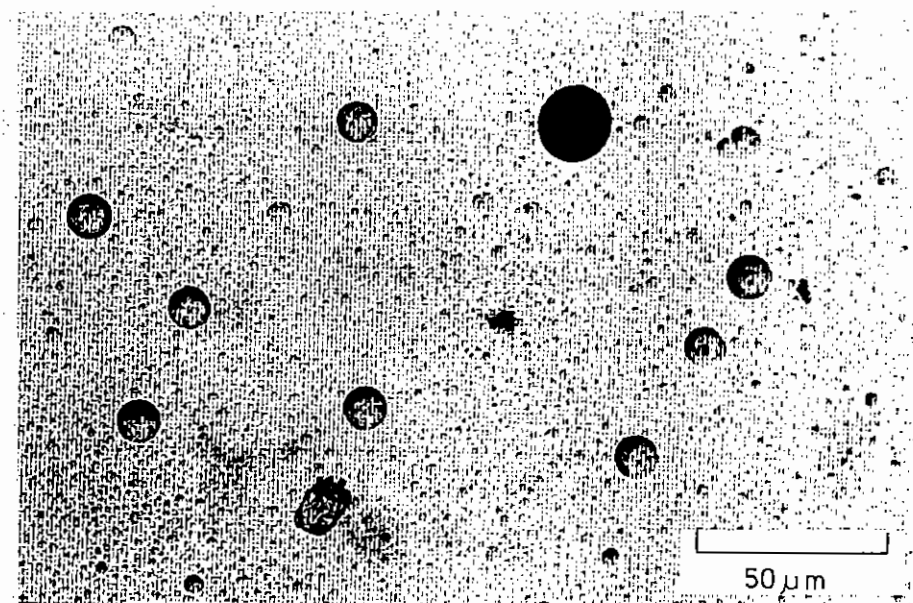


Fig. A2.42 A photomicrograph of α -particle tracks and a proton track from shot 14, 23/6/83.

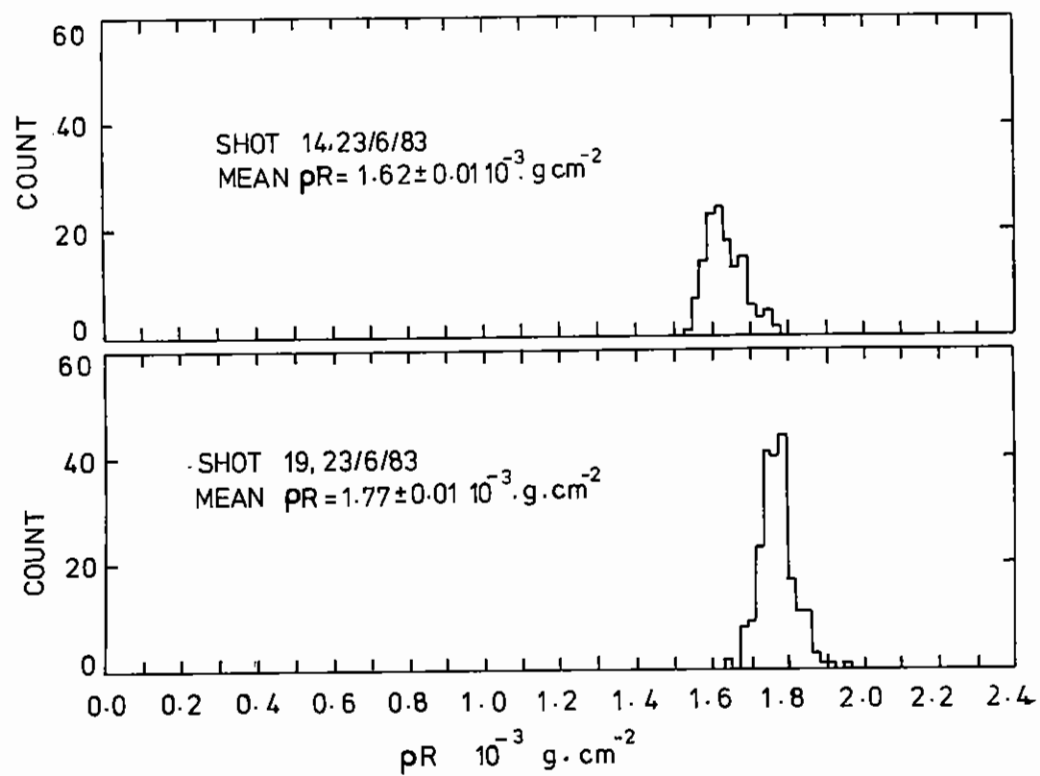


Fig. A2.43 Measured α -particle spectra from shots 14 and 19 on 23/6/83.

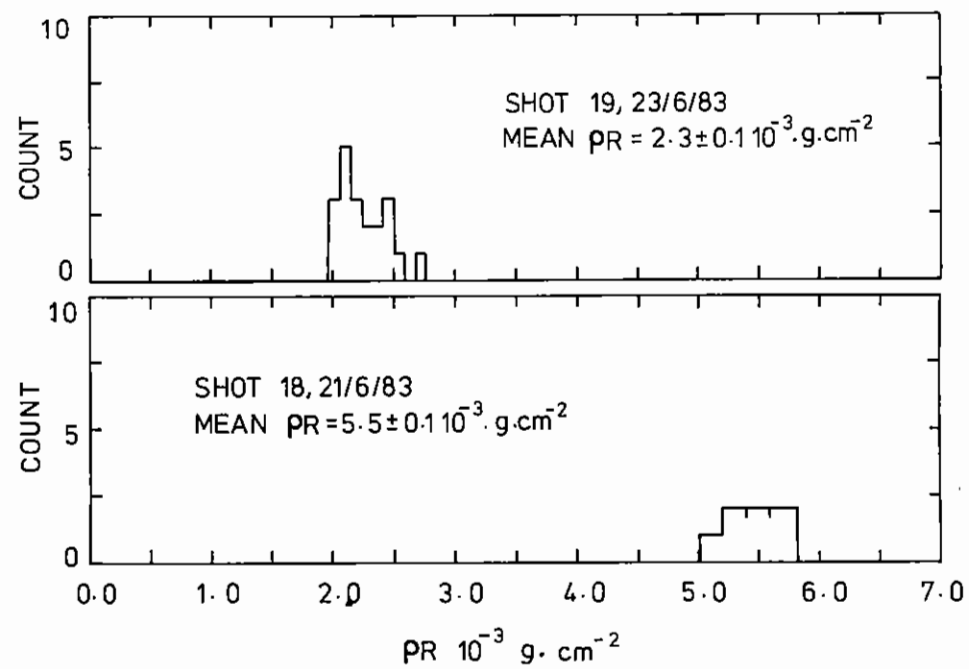


Fig. A2.44 Measured proton spectra from shot 19, 23/6/83 and shot 18, 21/6/83.

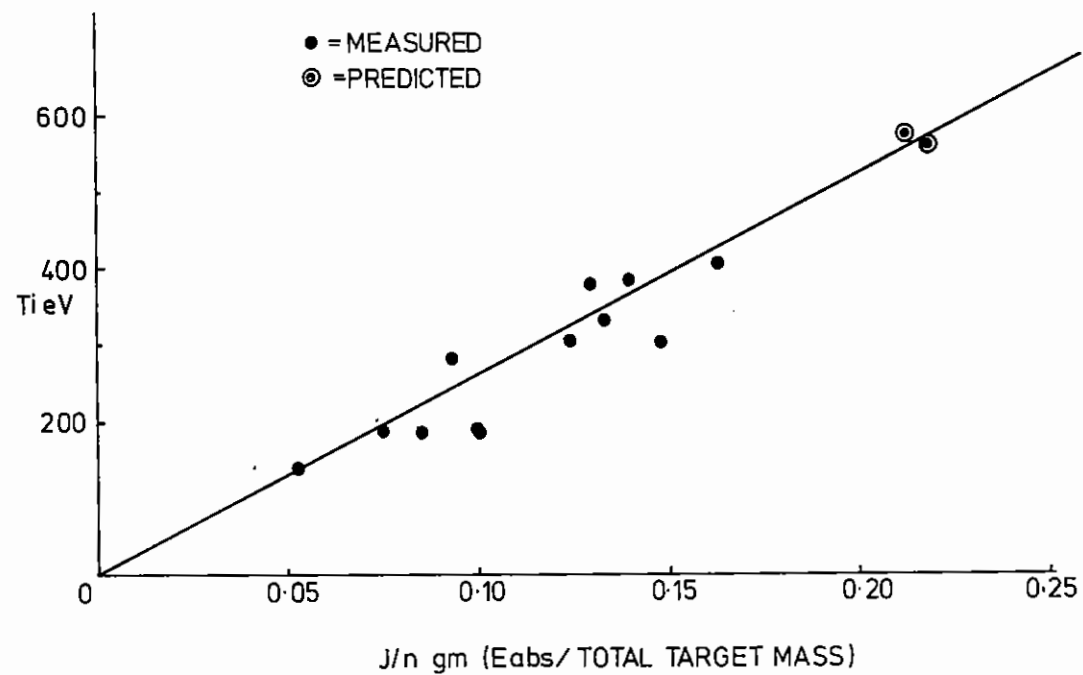


Fig. A2.45 Ion temperature deduced from neutron yield vs absorbed energy per unit mass.

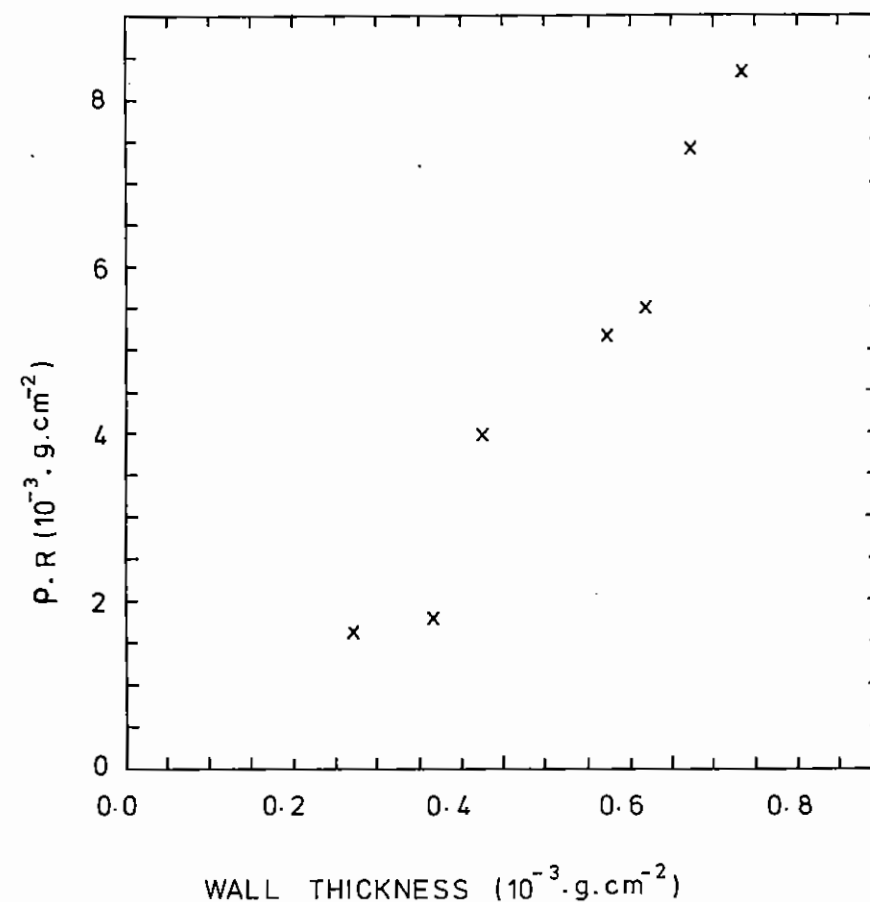


Fig. A2.46 A graph showing the variation of the density-radius product (pR) with initial microballoon wall thickness.

- (a) Straggling and instrumental broadening. This is estimated at $0.023 \times 10^{-3} \text{ g cm}^{-2}$ for alphas and between 0.13 and $0.27 \times 10^{-3} \text{ g cm}^{-2}$ for protons.
- (b) Doppler Broadening due to ion temperature. At a temperature of 375 eV the contribution to the sigma of the alpha particle distribution should be 0.028 gm cm^{-2} . The contribution to the proton sigma at the same T_i should be $0.23 \times 10^{-3} \text{ gm cm}^{-2}$.
- (c) Variation of $\rho\Delta R$ during the time of neutron emission.
- (d) Variation of $\rho\Delta R$ from place to place on the sphere.
- (e) Variations due to source size and source/shell aspect ratio effects.

These effects should all be combined in quadrature. The low values of $\sigma(\rho R)$ observed are consistent with the temperatures estimated on the basis of the yields. The residues, which must account for the combined effects of (c), (d) and (e) are so small that they impose tight constraints on the interpretation of the shots. Although the analysis of this data is still preliminary we may take, for example, shot 19-23-06-83 where the yield indicates $T_i \sim 380 \text{ eV}$ for a fuel compression of 1000 or 460 eV for a compression of 75. The residual alpha $\rho\Delta R$ is 1.65% of the mean $\rho\Delta R$ or $0.029 \text{ mgm cm}^{-2}$. The ρR of the uncompressed fuel in this shot is $0.035 \text{ mgm cm}^{-2}$ so that we must assume alpha emission comes from a small fraction of the compressed fuel. A small localised source is also required so that the spread in shell $\rho\Delta R$ due to particles traversing it in directions significantly away from radial can fit the data: a source to shell aspect ratio of less than $1/5$ to $1/10$ is required. The $\rho\Delta R$ observed is four times the value for the uncompressed shell so that the emission time must be of the order of a few picoseconds, unless it occurs at peak shell compression. The finite area scanned for tracks together with the effects of Coulomb scattering in the shell indicate that an area of shell equivalent to a circle $20 \mu\text{m}$ in diameter (at the initial shell position) was sampled.

Similar conclusions can be reached using the proton data from shot

18-21-06-83 although here the restriction on source size from variations in fuel $\rho\Delta R$ is less severe. But here the area of shell sampled was still larger, being equivalent to a $40 \mu\text{m}$ diameter circle at the initial shell position. A variation of less than three percent in the shell $\rho\Delta R$ is required by the data. Note that the shell $\rho\Delta R$ is here 7.5 times its uncompressed value.

The proton ($\rho\Delta R$) results are subject to large statistical uncertainties. They are however very close to the limit set by straggling and the source temperatures deduced from the yields. It may in fact be necessary to invoke a reduction in straggling due to energy loss taking place on free rather than bound electrons in the ionised shell. Such a reduction is to be expected due to the large number of soft collisions with free electrons which give enhanced energy loss but make little or no contribution to straggling.

In both the shots discussed above, the data are consistent with the thermonuclear particles being produced by a converging shock in a small volume of the fuel at an early time during the compression. A more complete analysis of this data is in progress.

A2.3 Point Source, Quasi-Monochromatic Radiography

C L S Lewis, J McGlinchey, S Saadat, R E Corbett (Queen's University of Belfast) R W Eason, C Hooker, D Bassett (RAL)

A2.3.1 Introduction

X-ray radiography has been established as a powerful technique in diagnosing ablatively imploded targets when the final compressed plasma is not hot enough to take advantage of spectral emission diagnostics. To date it has been customary to generate an extended X-ray emitting plasma to backlight an imploding target and to image this target via a pinhole or a Kirkpatrick-Baez microscope onto a detector which may be passive or active. With long-lived backlighter emission the radiograph may be streaked with 1-D spatial information. With short-lived synchronised backlighters a radiograph frame may be recorded with 2-D spatial information. There are

some disadvantages associated. Spectral discrimination in the probe radiation is limited to broad-band filtering with foils or sometimes reflectivity cut-offs if grazing incidence optics are employed. This complicates Abel inversion procedures applied to the spectrally integrated opacities measured from a radiograph and demands a knowledge of the relative spectral content of the probe radiation. Furthermore the requirement of an extended source sets a limit on the highest useful X-ray energies which can be generated for a given backlighter laser pulse energy and hence a limit on the maximum ρr value in a target which can be diagnosed. At present this limit is $\rho r \approx 2 \times 10^{-2} \text{ g cm}^{-2}$ for $\sim 2.5 \text{ keV}$ backlighting radiation (cf. Section A2.2.4 *ibid*).

This section describes preliminary studies into an alternative backlighting technique which promises to overcome these difficulties. The method relies on producing a bright plasma with spatial dimensions smaller than the desired spatial resolution in the object plane and to record radiographs by point projection.

A2.3.2 Point-Projection Assessment

Quasi-point sources were produced by irradiating small diameter wires ($\phi > 10 \mu\text{m}$) with the laser beam focussed at the end of the wire and incident perpendicular to the wire cylindrical axis. The effective size of the point source used is the plasma as seen looking along the wire axis from its free end. The simplest radiographic arrangement is to place an object on the wire axis between the wire end and a detector and to use broad band filtering for spectral selectivity. This is illustrated in Figure A2.47 where a copper grid is radiographed. Such radiographs can be used to estimate the scale size of the source in orthogonal directions and for $10 \mu\text{m}$ gold wires irradiated with 20J of energy in 100 psec at $1.06 \mu\text{m}$ using a $f/2.5$ lens ($\sim 50 \mu\text{m}$ diameter focal spot) the source size was typically found to be in the region $6\text{--}15 \mu\text{m}$. The smaller dimension is found to be in the direction of the incident laser beam.

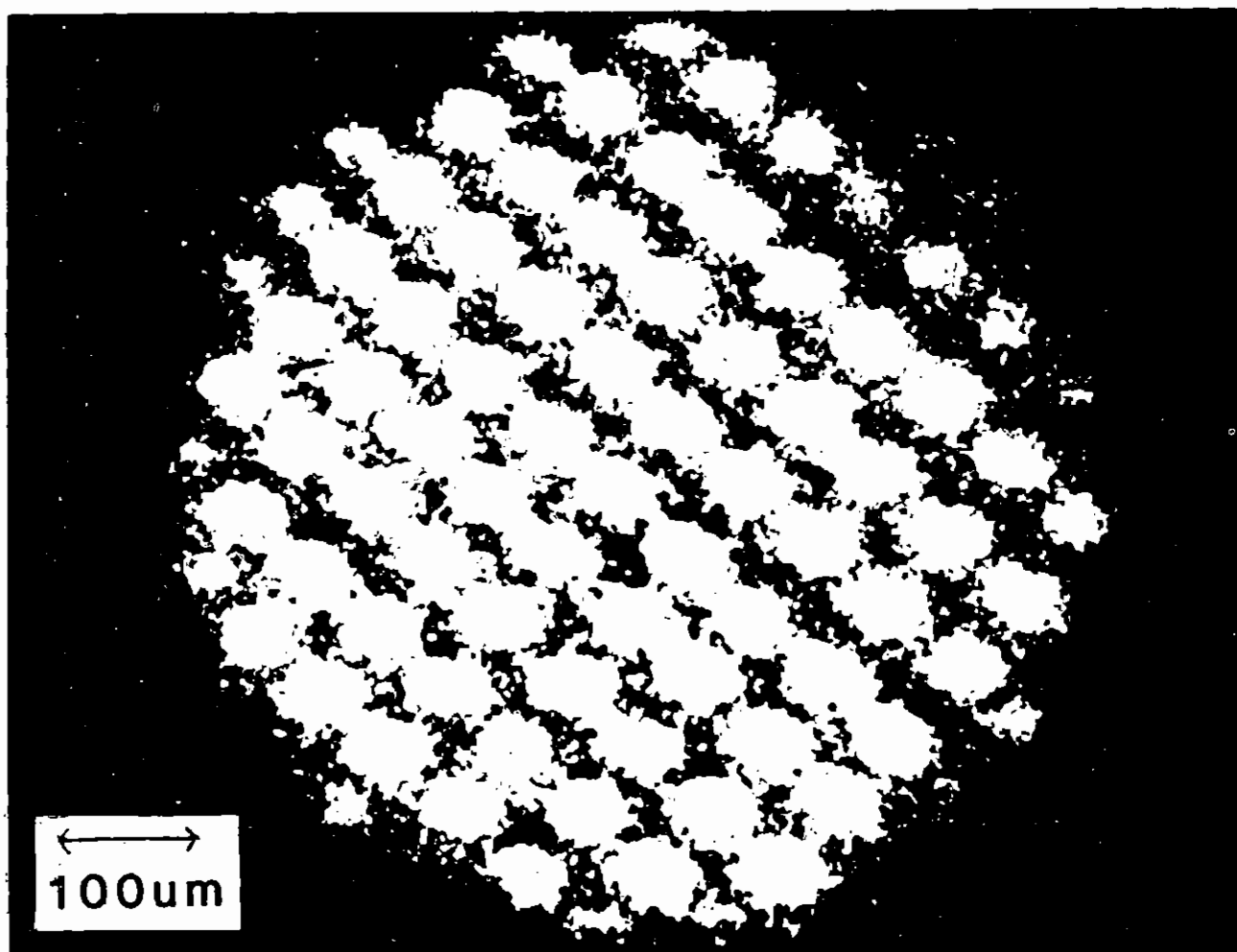
A development of the simple point projection is illustrated in Figure A2.48 where a Bragg crystal has been included to give spectral dispersion along

one of the spatial dimensions in the radiograph. This means that a wedge of rays from the source with photon energy satisfying a Bragg reflection condition will traverse a particular plane of the spherical object to be radiographed. In the circularly symmetric radiograph each chord perpendicular to the dispersion plane corresponds to monoenergetic probe radiation and there is spatial resolution along the chord. The orthogonal spatial dimension in the object is equivalent to spectrally displaced chords in the radiograph.

The geometry of the system studied is described with reference to the symbols in Figure A2.49 which shows the energy selective point projection system in plan and elevation.

An active X-ray detector was used and consisted of a P11 phosphor coated fibre optic plate coupled to an image intensifier outside the vacuum. The detector is illustrated in Figure A2.50. The phosphor plane was at the chamber wall at a distance p from the backlighter and had a useful diameter s . The target and crystal reflecting surface must be contained within the solid angle determined by the detector. i.e. $\Delta\Omega = (s \cos \theta/p)^2$ where θ is the Bragg angle for wavelengths incident normally on the detector. The target must be positioned at a distance from the backlight $> Dv/s$ giving a maximum magnification of the target of $m \approx s/D$. The dimensions of the crystal required depend on its position. We choose to position it midway between the backlight and detector so that $c = d$ and the mean wavelength reflected could be varied by traversing the crystal along the perpendicular bisector. In this case the crystal dimensions are given by $w = s/2$ and $L = s/(2 \sin \theta)$. The range of wavelengths $\Delta\lambda$ from the backlight needed to radiograph the entire target is determined by the range of Bragg angles $\Delta\theta = (s \cos \theta)/p$ subtended at the detector. For a crystal with plane spacing d this gives $\Delta\lambda = 2d(1 - (\lambda/2d)^2)s/p$ where $\lambda = 2d \sin \theta$. In the present case these parameters had values $p = 50 \text{ cm}$, $D \approx 100 \mu\text{m}$, $s = 50 \text{ mm}$, $L = 45 \text{ mm}$, $w = 22 \text{ mm}$, $u = 3 \text{ mm}$, $m \approx 200$ and the detector was normal to Bragg reflected rays for $\theta \approx 35^\circ$. The crystal used was PET with $2d = 8.74 \text{ \AA}$ giving $\Delta\lambda \sim 0.5 \text{ \AA}$ for radiation centred at $\lambda = 5 \text{ \AA}$ ($E \approx 2.5 \text{ keV}$) with $\theta = 35^\circ$.

Convenient backlighting targets were found to be of $Z \sim 80$ as these emit a



RADIOGRAPH OF Cu GRID WITH 25μm X RAY SOURCE

Fig. A2.47 Direct projection of Cu grid. Periodicity 70μm, hole diameter 35μm and W wire target $\approx$ 25μm.

MONOCHROMATIC PULSED RADIOGRAPHY

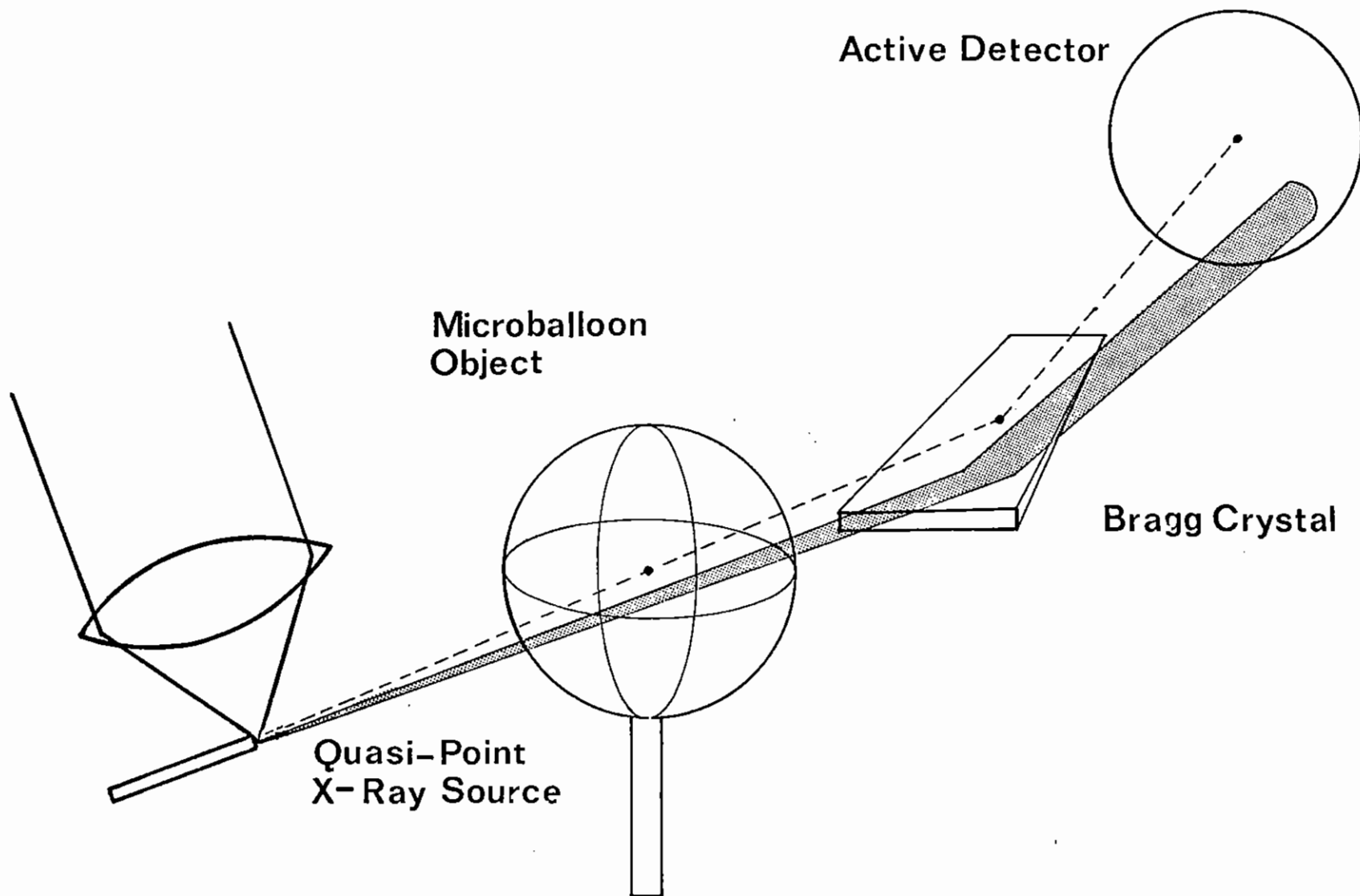


Fig. A2.48 Schematic of point projection with spectral dispersion.

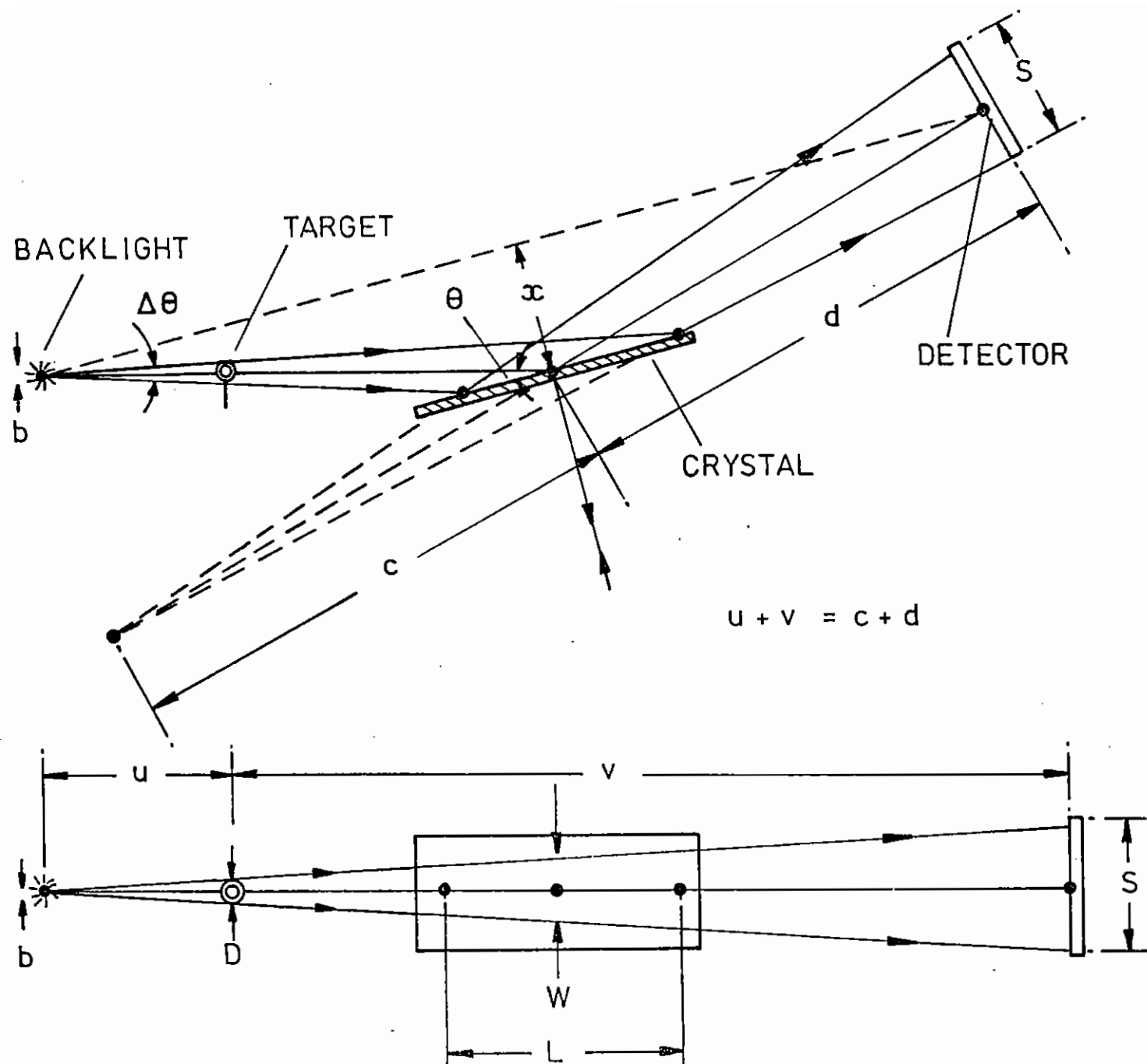


Fig. A2.49 Schematic of point-projection geometry with symbols used in text.

ACTIVE X-RAY DETECTOR

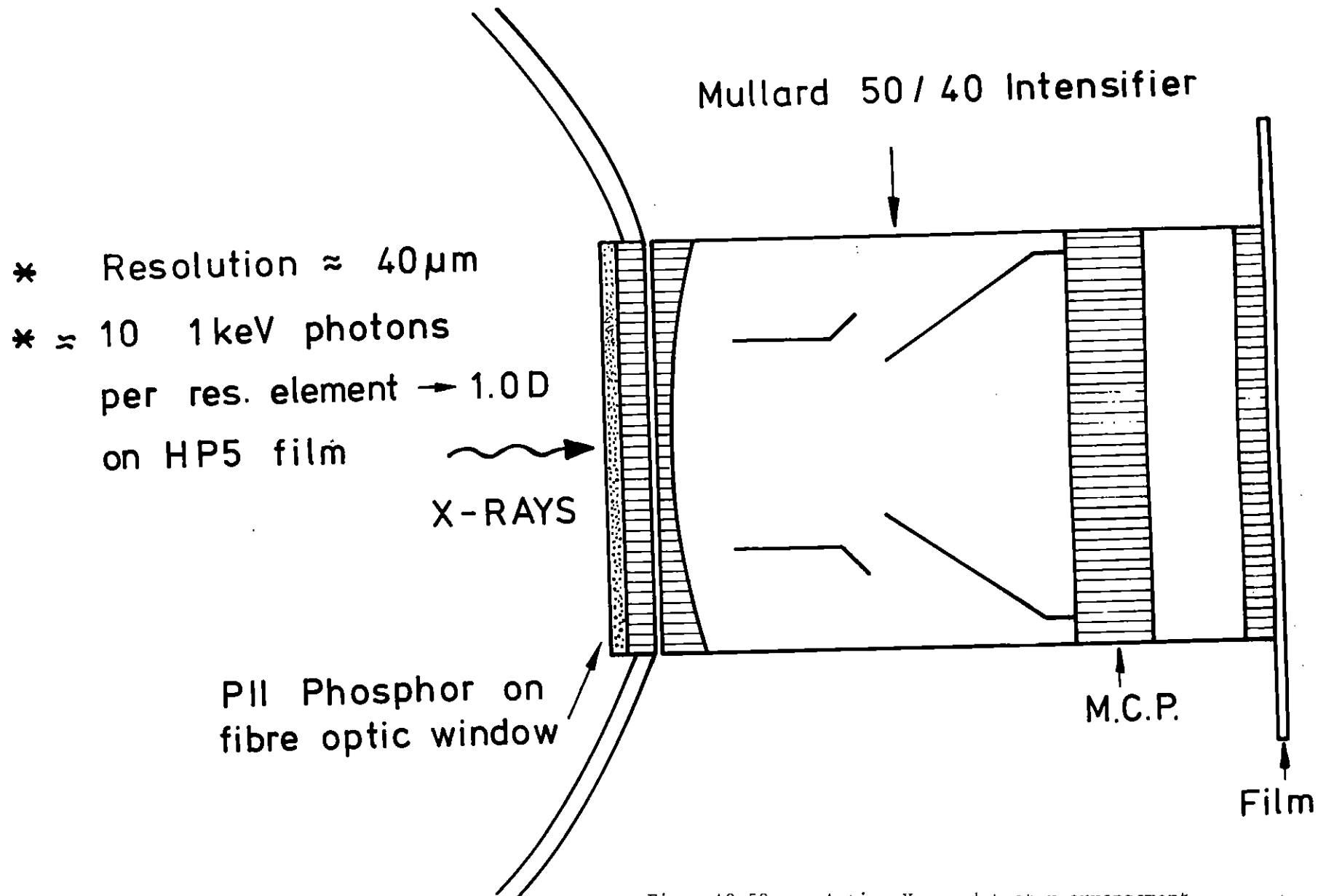


Fig. A2.50 Active X-ray detector arrangement.

bright quasi-continuous M-band spectrum with a suitable $\Delta\lambda/\lambda$. This is illustrated in Figure A2.51 where a range of targets were investigated using standard crystal spectrometers. The illustrated W spectrum is typical with the resolved lines in the region 4.2-4.8 Å attributed mainly to NaI-like 3-4 transitions and the quasi-continuum centred at ~ 6 Å attributed to CaI, KI, ScI, TiI, VI etc - like $3d^n - 3d^{n-1} 4p, 4f$ unresolved transition arrays.

The 10 μm wire backlightsers used were W, Au and Bi-coated gold.

A2.3.3 Experimental Observations

Spectra from W wires were recorded initially in the $\lambda = 4.2-4.8$ Å region to fix the crystal position in wavelength space. This was achieved by using a step wedge filter of 10 μm thicknesses of Saran foil ($\text{C}_2\text{H}_2\text{Cl}_2$), providing a wavelength marker at the Cl K-absorption edge at $\lambda = 4.4$ Å. Figure A2.52 illustrates a W spectrum with 5 Saran filter channels. Accurate centre wavelength setting could then be calculated as a function of the crystal translational parameter x in Figure A2.49. The step-wedge also allowed a relative exposure vs recorded film density response of the overall system. H-D characteristics were similar to the HP5 film's own HD curve implying constant conversion efficiency from X-rays to visible photons in the phosphor over the exposure ranges of interest. The narrowest line width observed showed that the spectral resolution of the system $\lambda/\Delta\lambda$ was ~ 1000 and that the integrated reflection coefficient $R \Delta\theta_c$ of the crystal $\sim 10^{-3}$ since spectral resolution was crystal rocking-curve limited.

Spherical target positioning for radiography was initially tested with solid Vanadium spheres of diameter $\sim 100\mu\text{m}$. Figure A2.53 illustrates such a radiograph using a gold backlight. The sharpness of the sphere edges indicate a spatial resolution of $\sim 10\mu\text{m}$. A background signal seen within the shadow is attributed to fluorescence from the crystal and inadequate light baffling of the phosphor from scattered light. The fluorescence can be drastically reduced by moving the crystal closer to the backlight and simultaneously reducing the exposed surface area. This however will make wavelength tuning more difficult. In retrospect the overall sensitivity of the system would allow much thicker Al filters in front of the phosphor to

A2.34

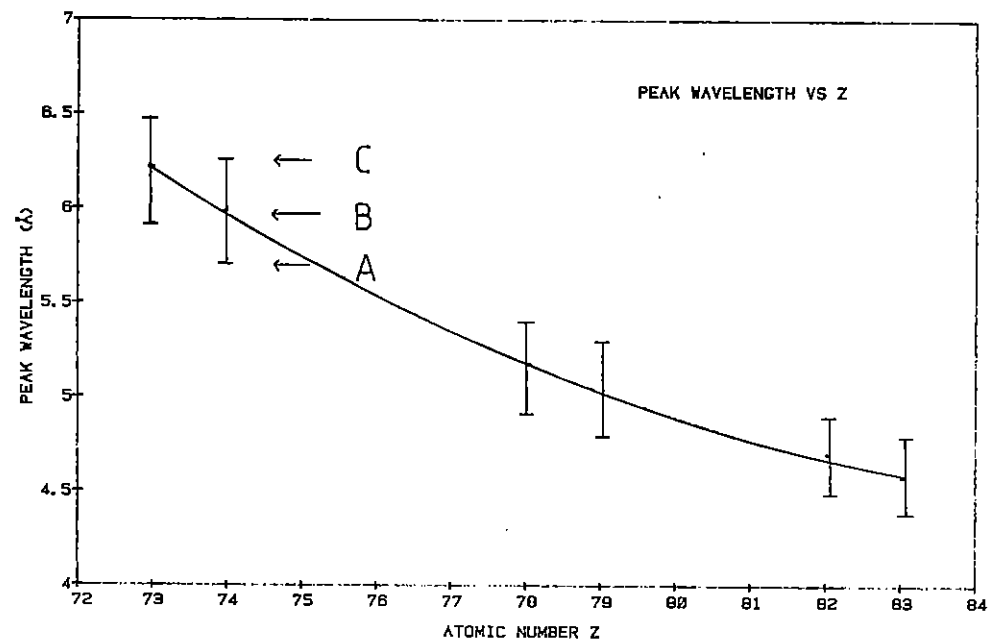
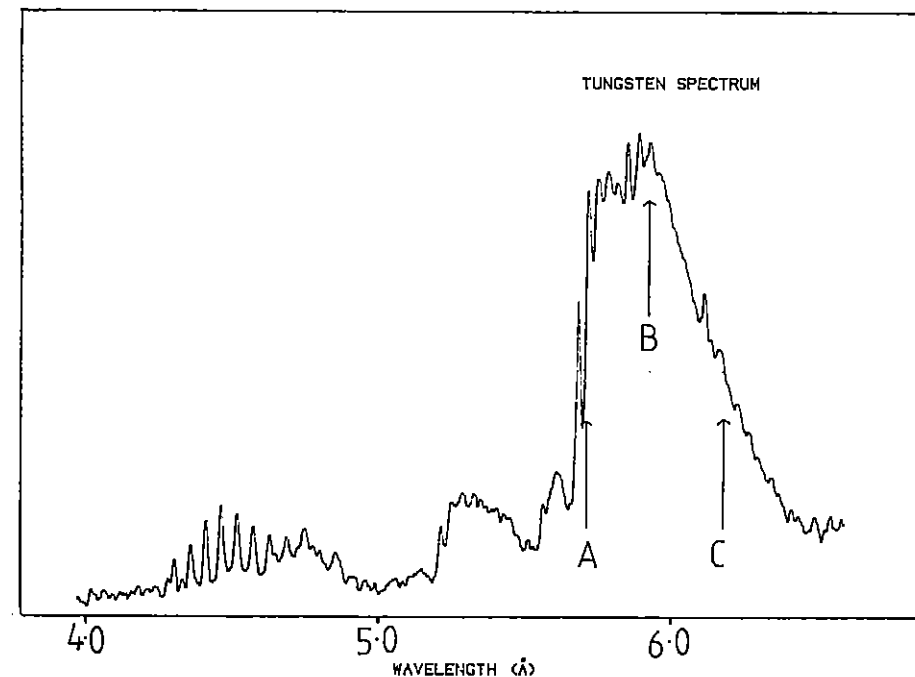
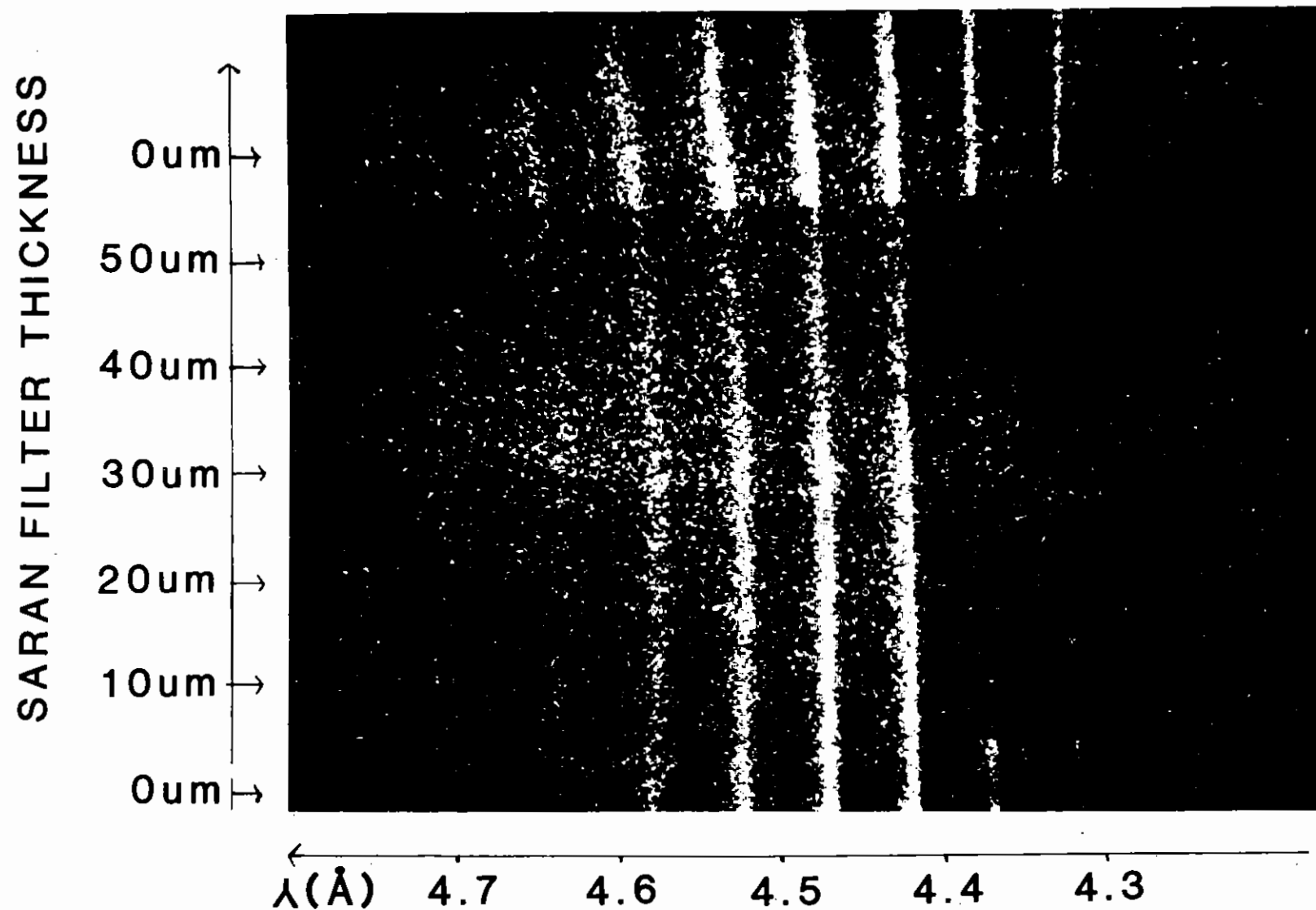
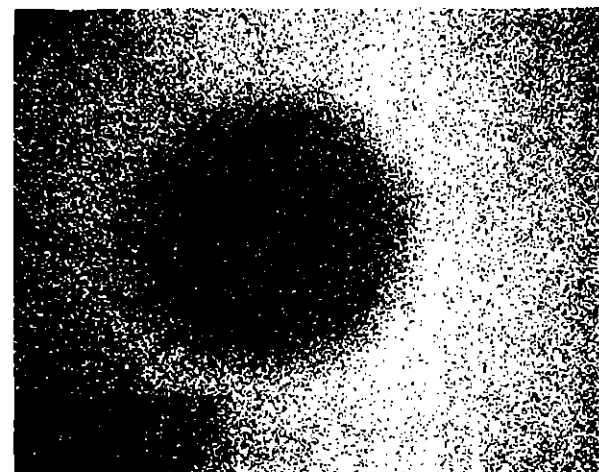


Fig. A2.51 Sources of useful spectra for point projection radiography.



21/10/83#3 TUNGSTEN SPECTRUM

Fig. A2.52 A tungsten line spectrum illustrating possibility of high resolution spectroscopy over a limited spectral range.



27/10/83 #7

GOLD B/L

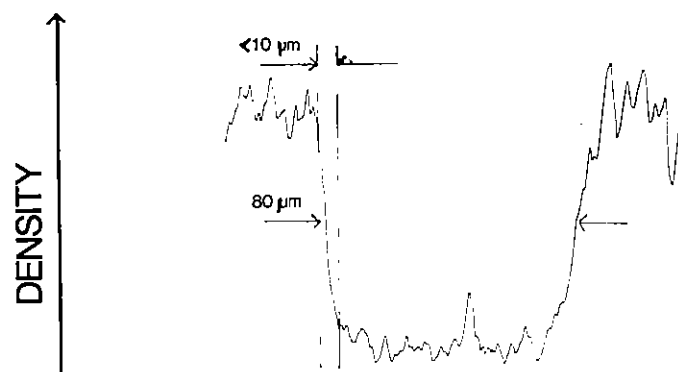
36J

100 ps

80 μ m

Vd SPHERE

λ (Å) 5.2 5.1 5.0 4.9 4.8



SHADOWING SOLID SPHERICAL TARGETS

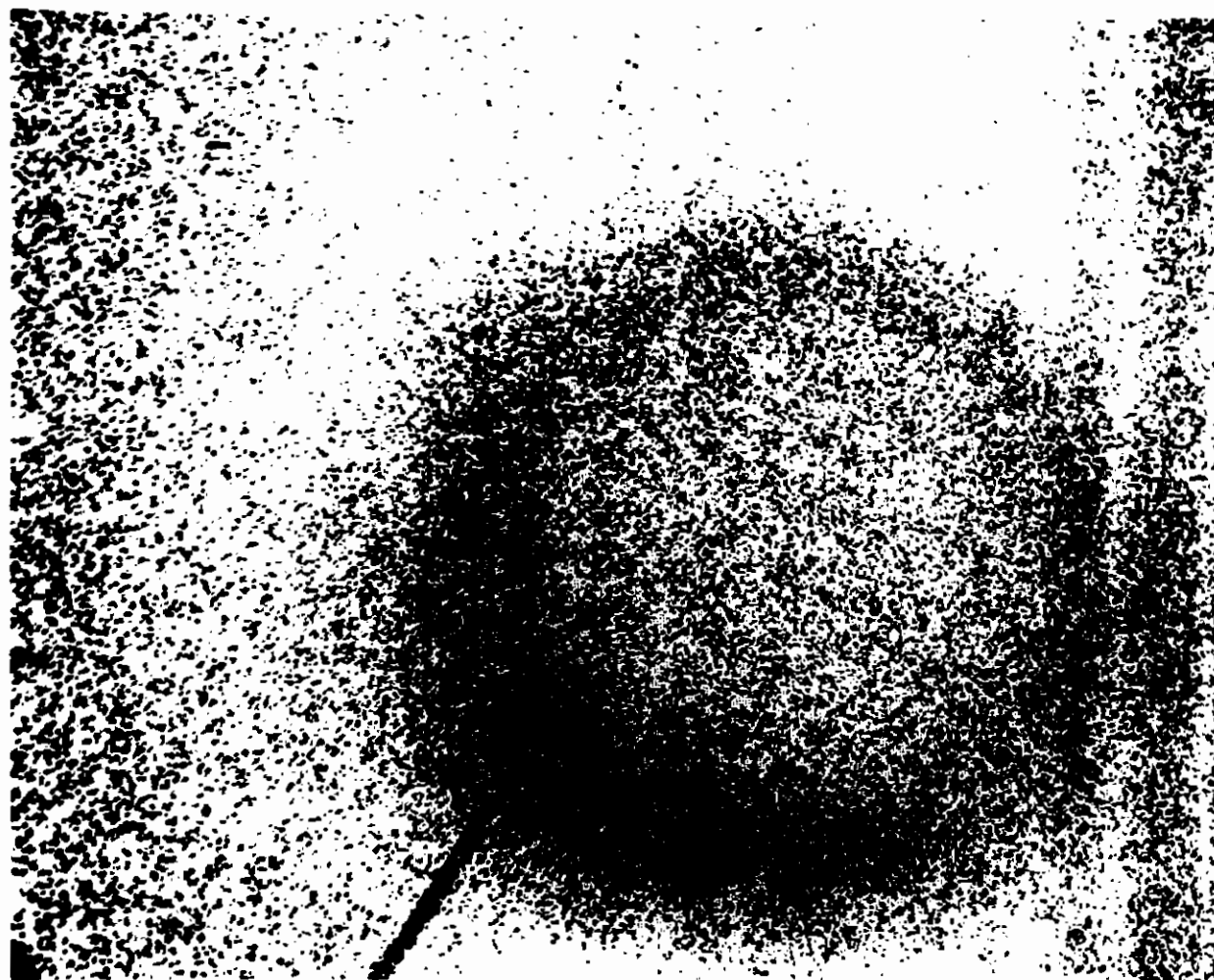
Fig. A2.53 Radiograph of a solid sphere.

suppress scattered light. In the experiment 6 μ m mylar with $\approx$ 700 Å Al on each surface was used.

Figure A2.54 illustrates an unirradiated CH shell target radiographed with a gold backlighter. The shell structure ($\Delta R \approx 9 \mu$ m) is easily discernable but the overall contrast is spoiled by the background discussed above. Finally we attempted some radiographs of imploding CH shells. Only 4 of the 6 Vulcan beams were available. These were symmetric but contained in a plane orthogonal to the radiography dispersion plane. Each beam delivered $\approx$ 20 J in $\approx$ 1 n sec at $\lambda \approx 0.52 \mu$ m and the radiograph was taken $\approx$ 600 psec after the peak of the drive pulse. A radiograph is shown in Figure A2.55 where the expected egg-shape of the resulting compressed core is seen. Self emission from irradiated targets does not appear to be a problem for low z CH targets due to the spectral selectivity of the crystal and the significantly greater brightness of the backlight in the selected spectral range. At worst the self-emission is a subtractable contribution in the radiograph.

A2.3.4 Conclusions

Point projection radiography with spectral dispersion has been demonstrated and with some technical improvements promises to be a powerful diagnostic technique for high pr measurements. Absolute brightness from the small diameter Au wires in the 2.5 keV region is $\sim$ 10-20% of the brightness of typical extended backlighters for the same incident energy and can be improved with tighter focussing onto the wire. This was estimated from pinhole images/projections onto X-ray film and from crystal spectrometers. The relative exposure density onto the detector in each case can be estimated by assuming, for the same spatial resolution, the imaging pinhole must have the same dimensions as the point source i.e. $b \sim 10 \mu$ m and that the extended source must be at least as big as the object to be radiographed i.e. $D \sim 100 \mu$ m. For a high magnification system $v \gg u$ the ratio of point/pinhole exposures is $\approx u \Delta \theta_c / d^2$ assuming the probe spectrum is flat and limited in each case to the wavelength range needed to radiograph an object of diameter D . In the present case, with 10-20% absolute emission from the backlighter the point projection image densities are comparable to



27/10/83 #1

GOLD B/L

31J

100 ps

POLYMER

$r=60\text{ }\mu\text{m}$ $r/dr=6.7$

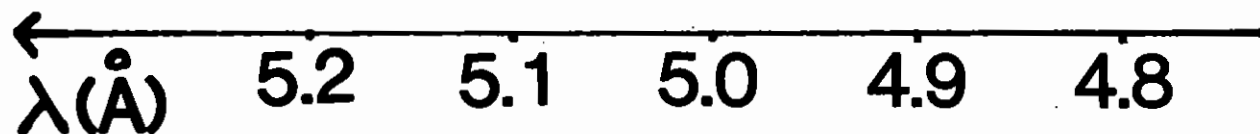
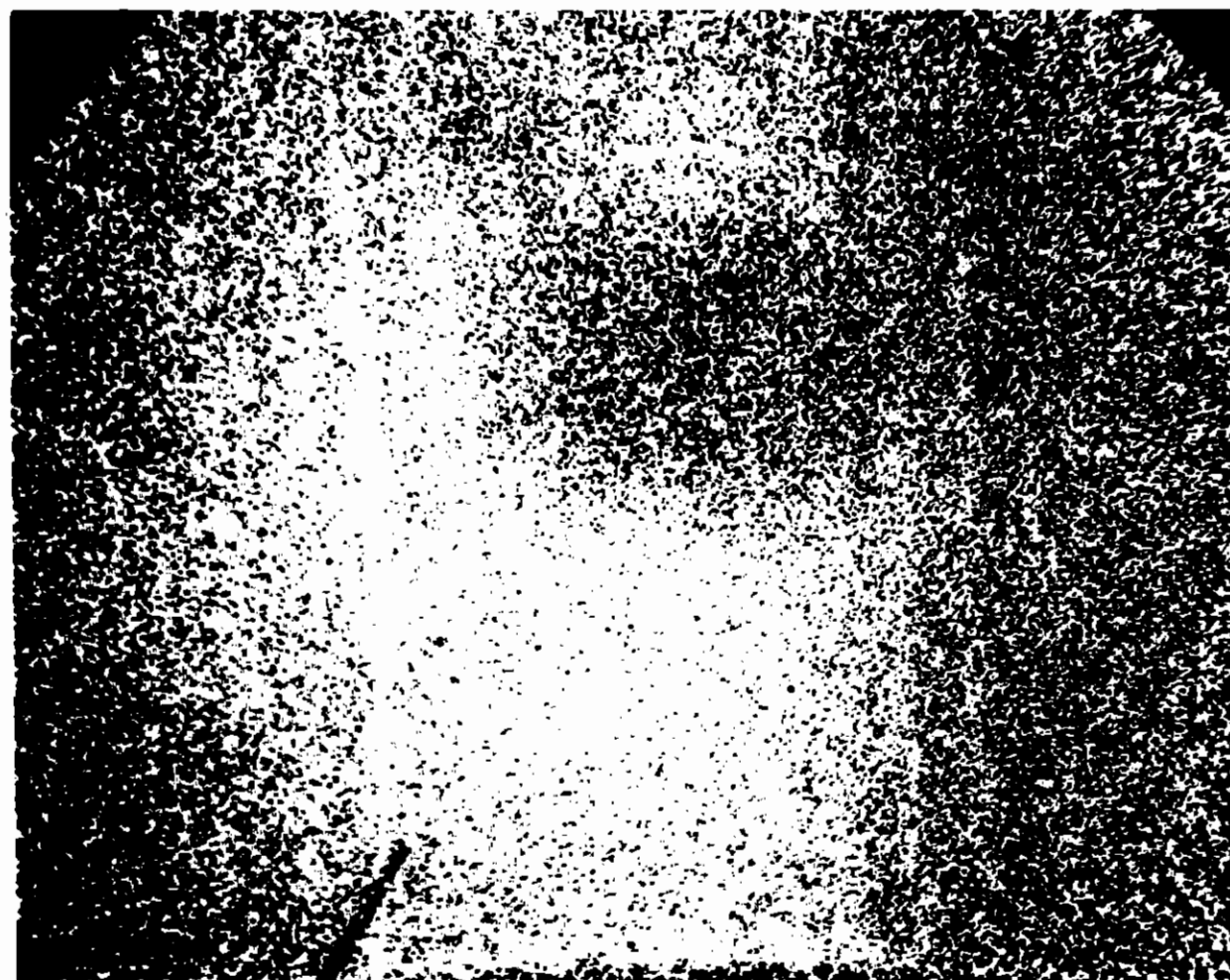


Fig. A2.54 Radiograph of a shell target.



27/10/83 #2

GOLD B/L

27J 100 ps

POLYMER

$r=68\text{ }\mu\text{m}$ $r/dr=6.7$

IRRADIATION

87J 4 BEAMS

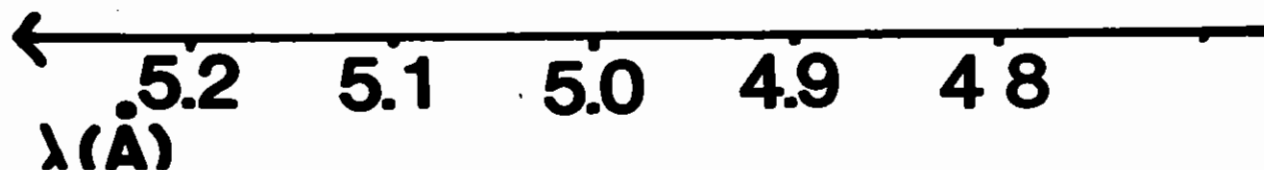


Fig. A2.55 Radiograph of an imploding shell target.

those obtained in pinhole imaged radiographs. This has been verified experimentally.

The point projection system has the potential to generate much harder probe radiation due to the higher incident laser power density and this may be in the form of spectral emission of strong resonance lines if 1-D spatial information on a single chord is adequate. Combinations of suitably placed resonance lines could generate a quasi 2-D radiograph. A further extension of the scheme is to replace the Bragg crystal with a narrow-band X-ray reflecting Barbee mirror. Here the probe radiation will contribute to all parts of the image but can be made quasi-monochromatic with an optimally designed mirror for broad continuum sources or by choosing a suitably strong resonance line within the mirror pass-band.

A2.4 Framing Camera Studies of 1.05 μ m Implosions

N Finn, T A Hall and E McGoldrick (Essex)

Development work on the x-ray framing camera (see section A6.7.1 and reference A2.21) has proceeded far enough for it to be possible to use it as a diagnostic on the six beam system.

The camera was used to observe the collapse of microballons illuminated by six beams of 100J, 1.05 μ m light in a 2ns pulse, using standard backlighting techniques. The balloons were of glass shells of 130 - 180 μ m diameter with 5-6 μ m thick plastic coatings to reduce self emission of X-rays while maintaining high absorption of backlighting X-rays at the core. Some of these were etched with 40 μ m etchings on one side (A2.22). The backlighting target was copper irradiated by a 50J, 2ns, 1.05 μ m laser pulse focussed to a 250 μ m spot. The defocus of the six beams was kept at a D/R of 4.5, and the backlighting and main beams were co-incident in time.

The Si switch of the framing camera was switched with a series of short (~100ps) pulses produced by a Fabry Perot etalon. To produce pulses of equal intensity (to give equally spaced frames) the etalon was slightly mis-aligned to separate the pulses in space. The beams then passed through

a neutral density filter wedge to produce equal intensity pulses. These were then focussed by a lens onto a diffusing screen in front of the switch.

A 10 μ m pinhole was placed ~2cm from the target to give ~ x 10 magnification at the Caesium Iodide photocathode, giving an overall image magnification on film of ~ x 30. The framing rate was set by the etalon at two frames per nanosecond (Figure A2.56).

Figures A2.57 and A2.58 are typical of the results obtained.

Four images can be seen (although the first in Figure A2.57 is very faint), separated in time by 500 ps, and the progress of the implosion can be clearly followed. The implosion is not spherical and takes on a hexagonal shape as it implodes due to the increased drive pressure on the balloon's surface where two and three beams overlap. Similar structure has been observed using a short back-lighting pulse and a pinhole camera (see section A2.2.4).

Detailed analysis of the data is being carried out, including the comparison of implosions with etched and unetched targets.

A2.5 Measurements of the Rayleigh-Taylor Instability in Laser Irradiated Targets

J Wark, J D Kilkenny (IC), A J Cole, P T Rumsby and M H Key (RAL)

Recent measurements of the Rayleigh-Taylor (R-T) instability (A2.23) in laser accelerated plane targets show that the growth rate, γ , is approximately one third of the classical value (A2.1), $\sqrt{(ka)}$, where k is the wavenumber and a the acceleration. The experiments discussed here show the first clear evidence for spike formation in a laser driven R-T instability.

These results were obtained using a target with an initial mass perturbation of known k . The targets were fabricated by a novel laser-etching technique as described in reference (A2.22).

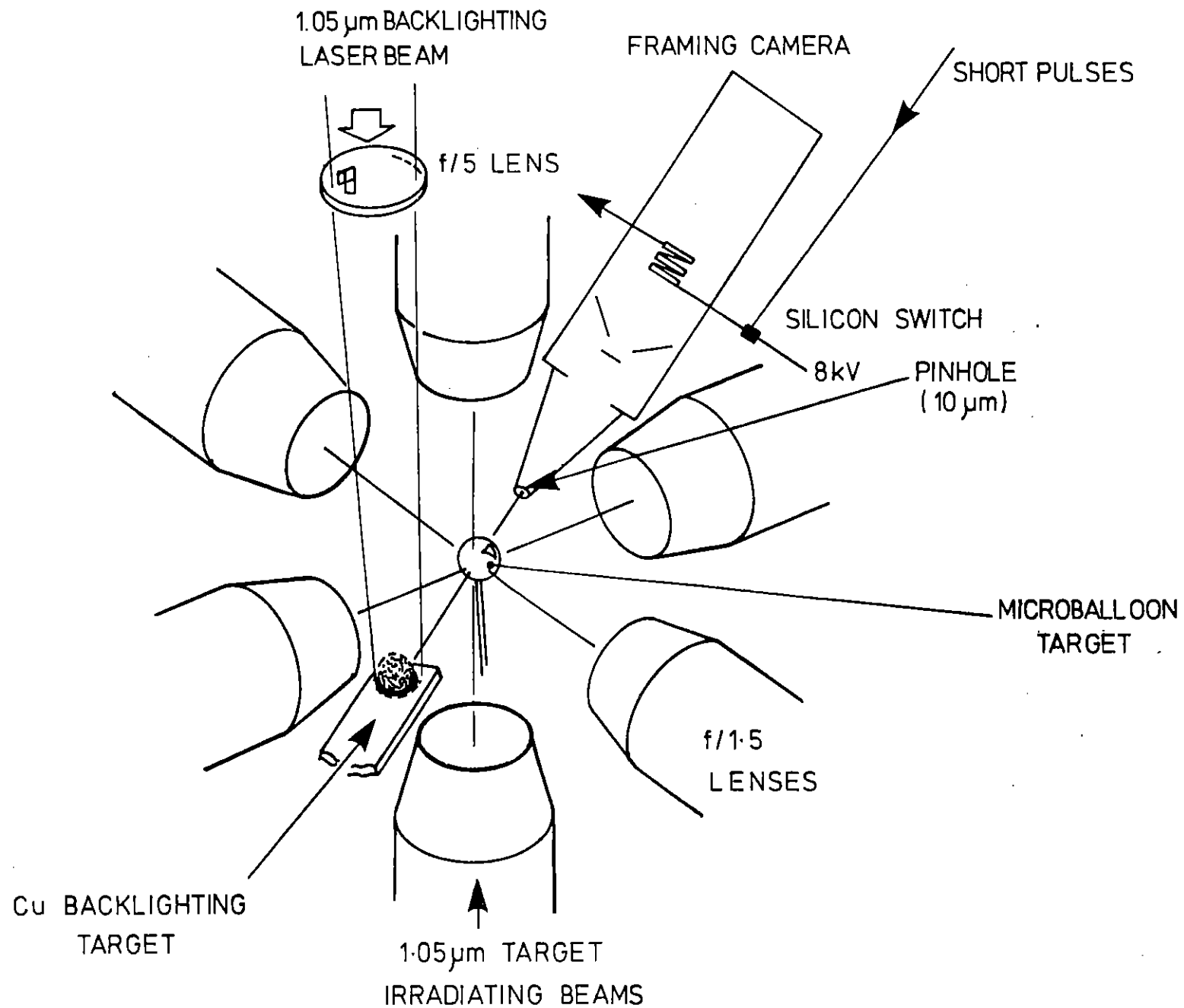


Fig. A2.56 Schematic diagram showing the position of the framing camera and backlighting source relative to the microballoon target.

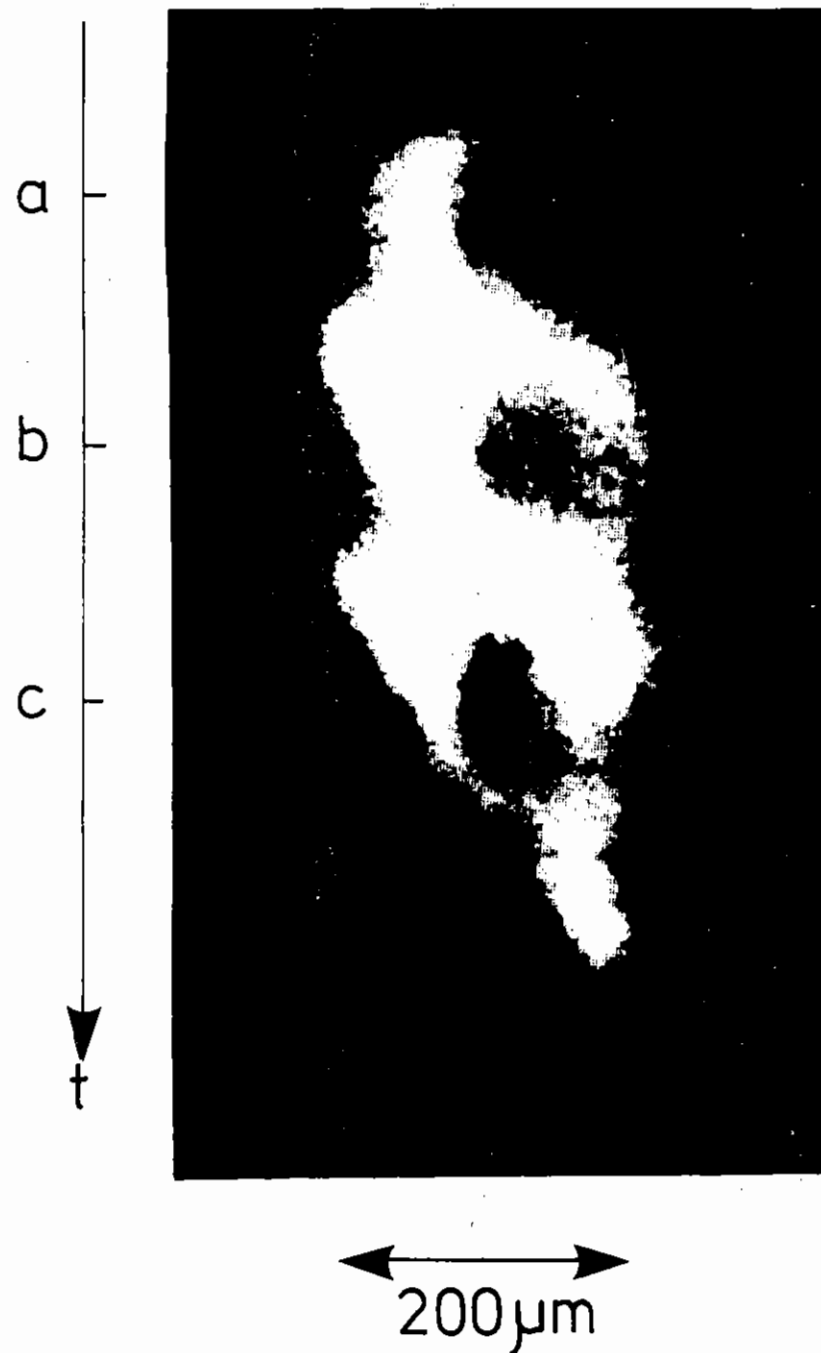


Fig. A2.57 Framing camera pictures of a 180μm diameter plastic coated glass sphere imploding.

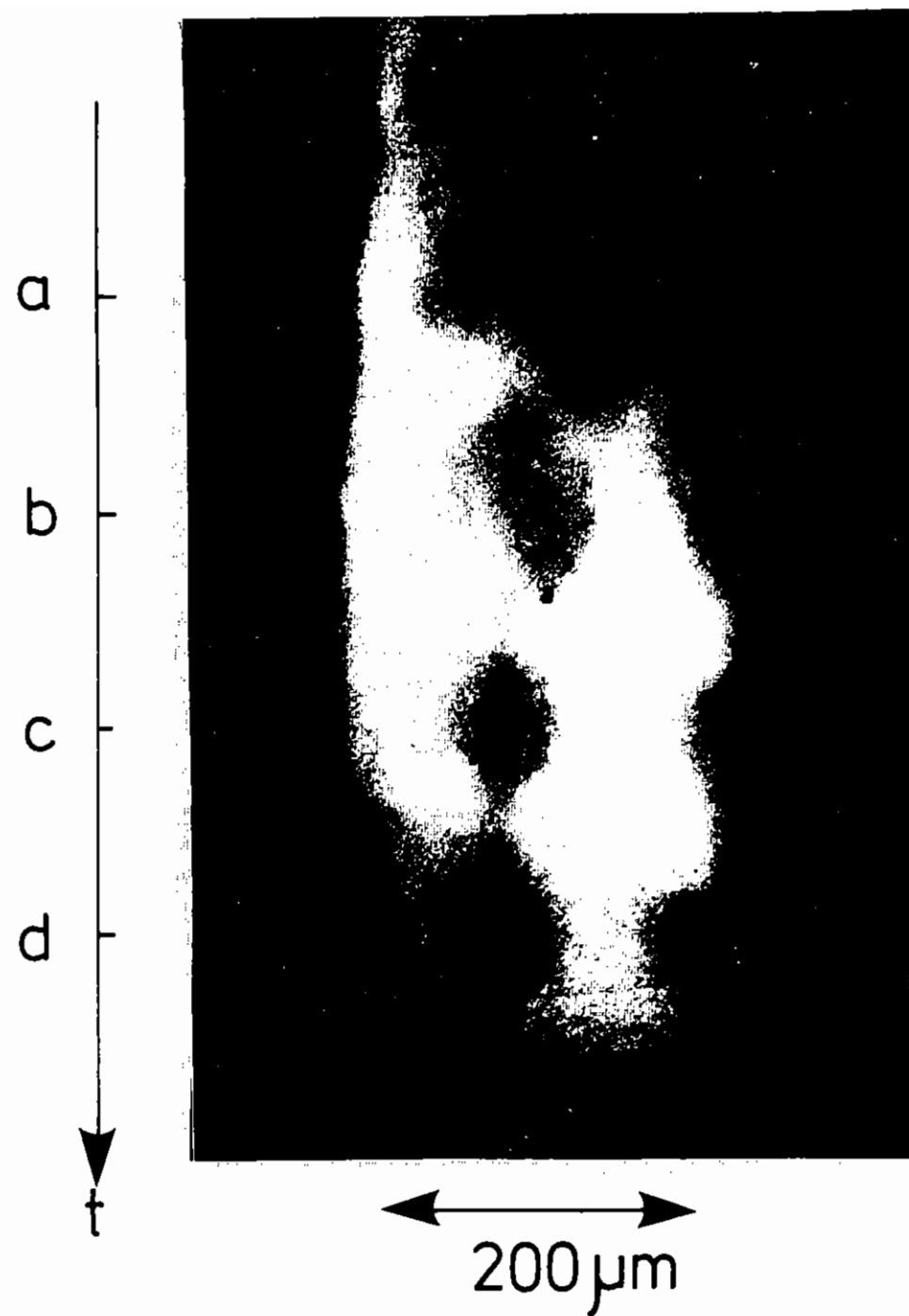


Fig. A2.58 Framing camera pictures of a 165μm diameter plastic coated glass sphere imploding.

The targets were plastic microballoons of $\sim 150 \mu\text{m}$ radius and $10 \mu\text{m}$ wall thickness. One half of the target was mass modulated with a wavelength of $20 \mu\text{m}$ and a modulation depth of $\sim 2 \mu\text{m}$ as shown in Figure A2.59.

The experimental set-up is shown in Fig A2.60. The etched side of the targets were irradiated with three of the six orthogonal $0.53 \mu\text{m}$ beams of the six beam irradiation facility. The beams were focussed 4.5 target radii beyond the spheres for optimum uniformity of illumination (see section A2.2.3 *ibid*) with a typical total absorbed irradiance of $4 \times 10^{13} \text{W cm}^{-2}$. Behind the microballoon a separate Cu target was irradiated by a fourth synchronous short pulse backlighting laser beam. This produced a plasma typically $300 \mu\text{m}$ across and $200 \mu\text{m}$ high. In the spectral region of observation the X-ray emission from the microballoon was negligible compared to the back-light. The Cu plasma provided a source with which to make an X-ray absorption picture of the irradiated microballoon. The absorption picture was recorded using an active pin-hole camera with a magnification of 20 and pinhole diameter of $9 \mu\text{m}$ and $3 \mu\text{m}$ Al filter. The spectrum of the Cu backlighting target, measured with a crystal spectrometer, showed a peak at a wavelength of $10 \text{\AA}$, with a rapid fall at shorter wavelength. The spectrum gives agreement with the transmission observed at the centre of the microballoon, calculated by extrapolation of the backlight intensity from outside the region shadowed by the target.

The delay of the backlighting beam with respect to the main beams was measured using an X-ray streak camera. This yielded the length of the X-ray pulse from the Cu plasma, which was found to be 300ps FWHM. The velocity of the microballoon wall is $\sqrt{2as}$; where s is the distance moved. Using independently measured pressure-irradiance scaling relationships, (A2.4), the maximum time-smearing resolution was calculated to be $\sim 8 \mu\text{m}$, and was typically less than the pinhole resolution.

An X-ray shadowgraph of the transmission of such a target is shown in Figure A2.61. In this case the irradiated side of the target has travelled inwards $15 \mu\text{m}$ and spikes protruding radially outwards from the sphere can clearly be seen. The unirradiated side of the balloon was used as a reference point to calculate the distance travelled by the imploding hemisphere. The amplitude

of the spikes cannot be explained by secular growth, i.e. the troughs of the initial perturbation moving further due to their smaller mass.

Assuming that the ablated mass is negligible (a maximum of 15% of the wall thickness in this experiment) then the classical growth exponent, γ_t , can be expressed as $\sqrt{2sk}$, where S is the distance moved by the target. The average amplitude of the observed spike is plotted as a function of $\sqrt{S}$ in Figure A2.62. Also shown is the expected amplitude due to purely secular growth allowing for a typical mass ablation rate calculated from independent measurements (A2.4). The total growth rate is $0.6 \gamma_c$; subtracting the secular growth gives $0.25 \gamma_c$.

Subsequently another two microballoons were shot that had an initial perpendicular amplitude of only $1 \mu\text{m}$. The results are shown in Figure A2.63. A modulation of the correct wavelength can clearly be seen in Figures A2.63(a). The etched hemisphere of the second microballoon was accelerated towards the pinhole camera, rather than perpendicular to it, with the unmodulated side thinned down to $0.5 \mu\text{m}$ Figure A2.63(b). A horizontal bar can clearly be seen through the thinned part of the balloon showing that mass modulation has indeed taken place.

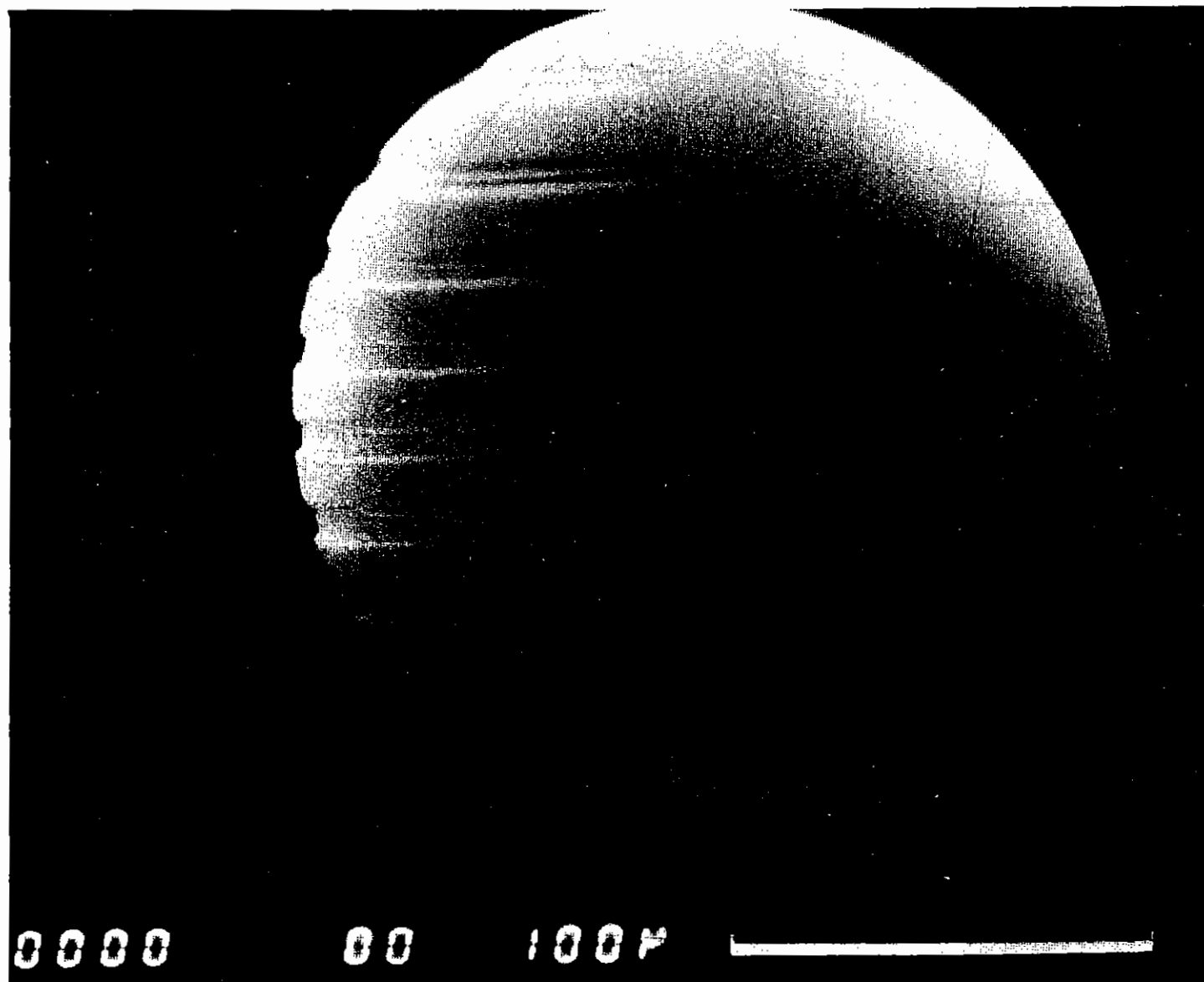
In conclusion we have observed the growth of radially protruding spikes on an irradiated sphere with an initial mass perturbation. The amplitude of the spikes cannot be explained purely by secular growth, but agree with a Rayleigh-Taylor growth rate of ~ 0.25 of the classical. Targets with a lower amplitude initial perturbation also show evidence of spike formation. Acceleration of such a target towards the pinhole camera has shown clear evidence of mass redistribution.

A2.6 Thermal Smoothing Measurements at $0.35 \mu\text{m}$

A J Rankin, J D Kilkenny, S D Tabatabaei (IC) and A J Cole (RAL)

A2.6.1 Introduction

Recent experiments at this laboratory (A2.24) have demonstrated the use of a shock transit time technique to measure the degree of smoothing in non



modulation wavelength : 20 μ m

Fig. A2.59 Electron micrograph of an etched polymer microballoon target.

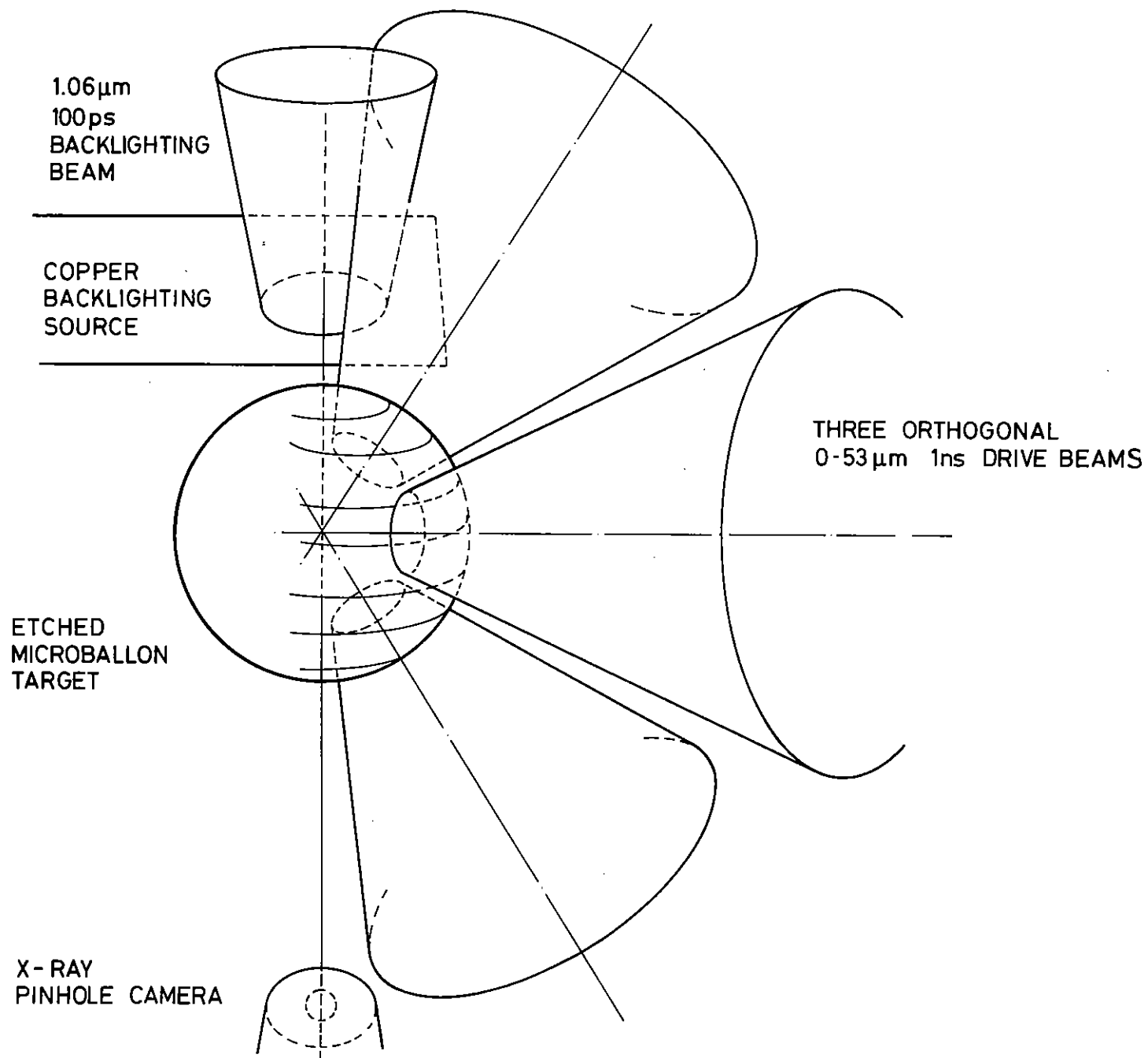
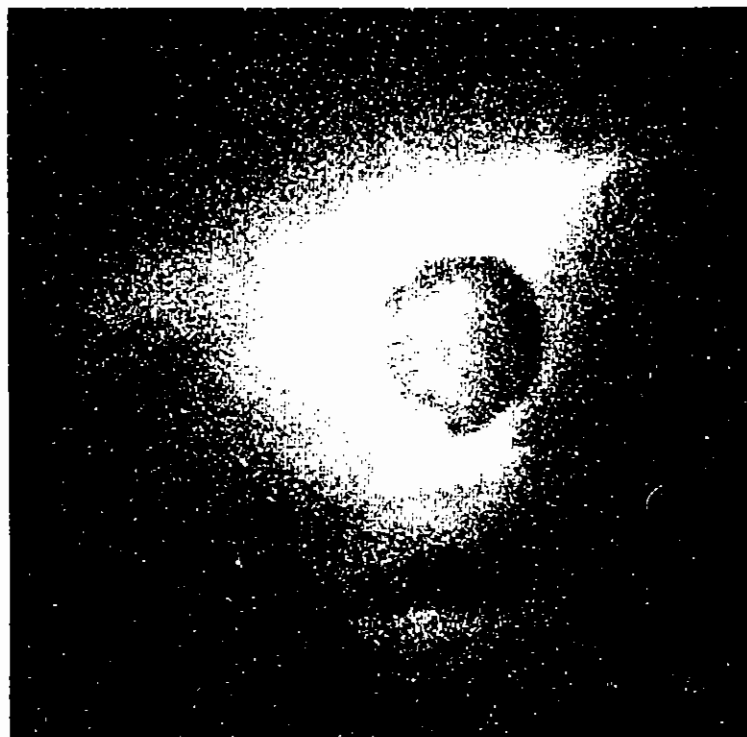


Fig. A2.60 Experimental arrangement for one-sided irradiation of etched microballoons.

laser



shot : 10280783

Ro : 91 μm

Ia : $3.7 \cdot 10^{13}$

delay : 0.75ns

S : 14.4 μm

Fig. A2.61 X-ray radiograph of an irradiated etched microballoon showing spike growth.

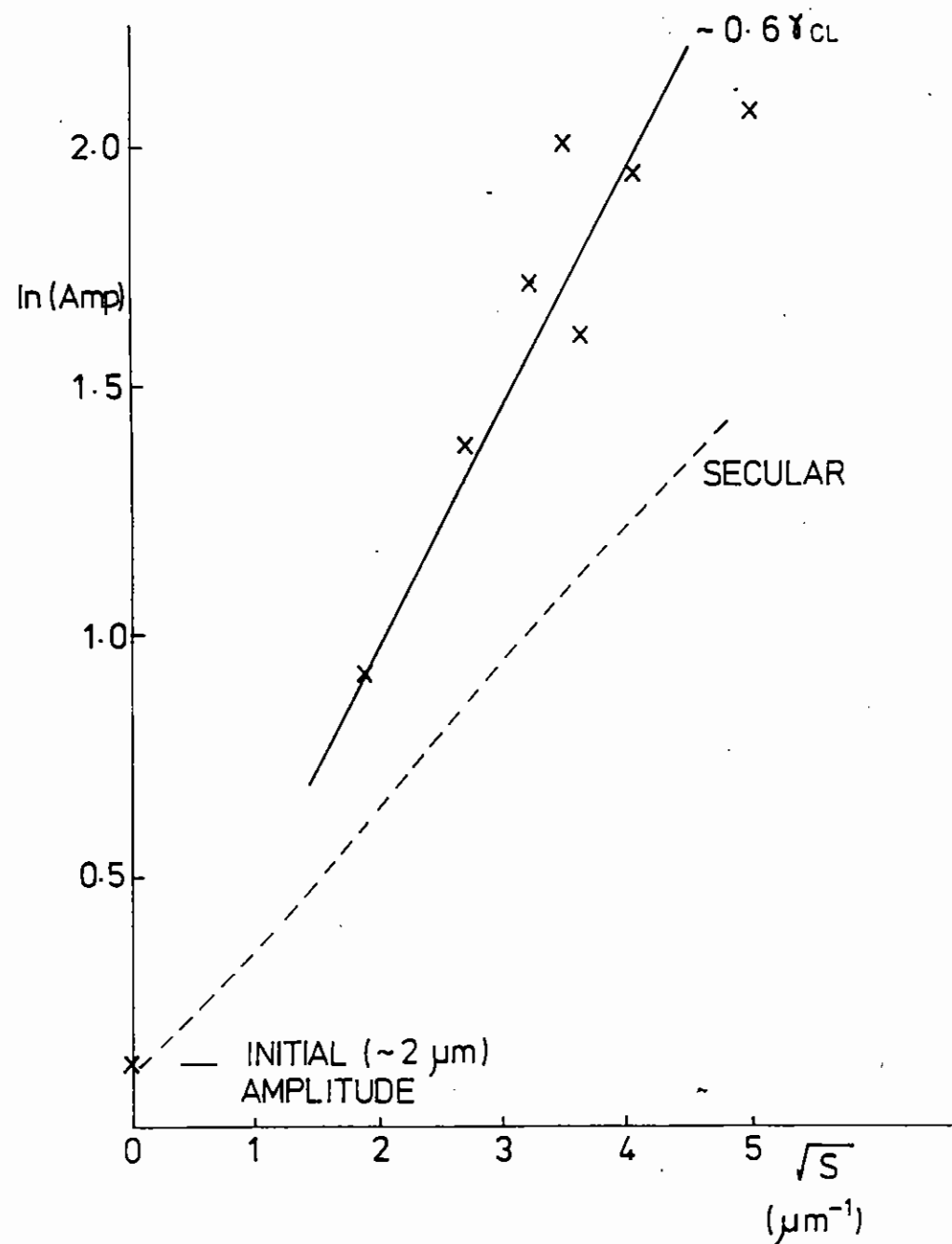


Fig. A2.62 Plot of amplitude of the spikes against distance moved by the irradiated side of the target. The calculated secular growth is shown for comparison.

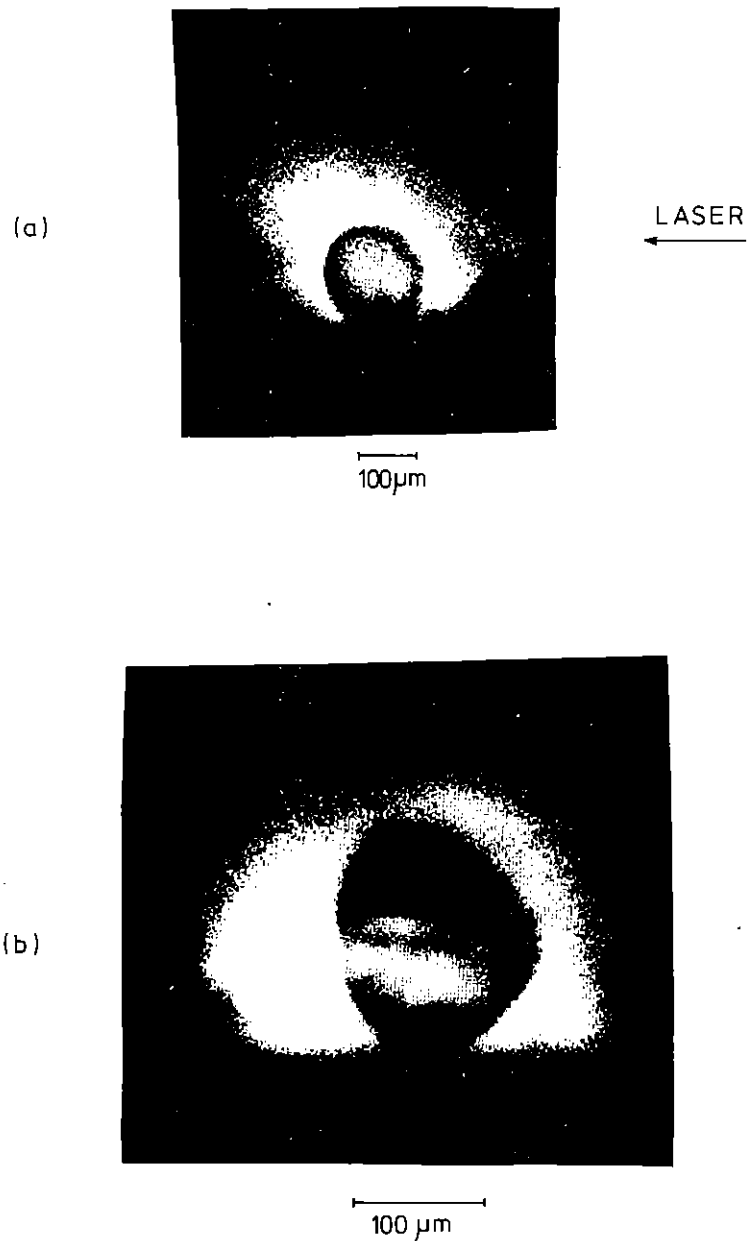


Fig. A2.63(a) Side-on radiograph of an etched balloon with an initial perturbation amplitude of $1\mu\text{m}$.

(b) Face-on radiograph through the thinned region on the unirradiated side of the target. Horizontal bars of increased absorption are visible.

uniformly irradiated foil targets. The results obtained previously suffered from an anomaly concerning the zero smoothing shock transit time ratio and the very limited irradiance data set obtained. This second experiment has succeeded in resolving this anomaly and expanded the irradiance regime considered to reveal the anticipated smaller degree of smoothing from a technique avoiding hydrodynamic foil acceleration effects (A2.9). Improvements in image quality have been made through more efficient relay optics and evidence of filamentation or jet structures on pinhole camera photographs also obtained.

A2.6.2 Experiment

Figure A2.64 shows the experimental set up. A single beam of the Vulcan Nd glass laser was frequency tripled in KDP crystals to deliver up to 30 J in a 600ps pulse. The focal spot sizes were typically $150\mu\text{m}$ and were monitored through a pinhole camera with $10\mu\text{m}$ Pt pinhole, viewing side on to the $25\mu\text{m}$ thick aluminium targets. The imaging system used an f1.2 50mm camera lens equipped with a light exclusion cone to relay the image through a specially designed f1.6 157cm achromatic lens onto the $250\mu\text{m}$ wide camera slit.

$3\omega_0$ rejection mirrors protected the Hadland streak camera equipped with an S20 photocathode from scattered laser light. A laser beam timing fiducial was supplied as in previous experiments (A2.24) and beam modulations were again simple stripes of neutral density filter.

Figure A2.65 shows a typical optical streak image from a $25\mu\text{m}$ Al target for an unmodulated and modulated beam. Figure A2.66 shows equivalent plane microdensitometer tracings corresponding to the unmodulated and modulated beam cases.

The scaling of the ratio of the maximum and minimum transit time $(\frac{T_{\text{max}}}{T_{\text{min}}}-1)$ as a function of incident irradiance is shown in Figure A2.67. If no smoothing were present the drive pressure should scale simply as the inverse square, if the pressure scales as $P \propto I$ (A2.4) then using the measure irradiance

SMOOTHING MEASUREMENTS OF IMPOSED BEAM MODULATIONS USING A MODULATED SHOCK BREAKOUT TECHNIQUE

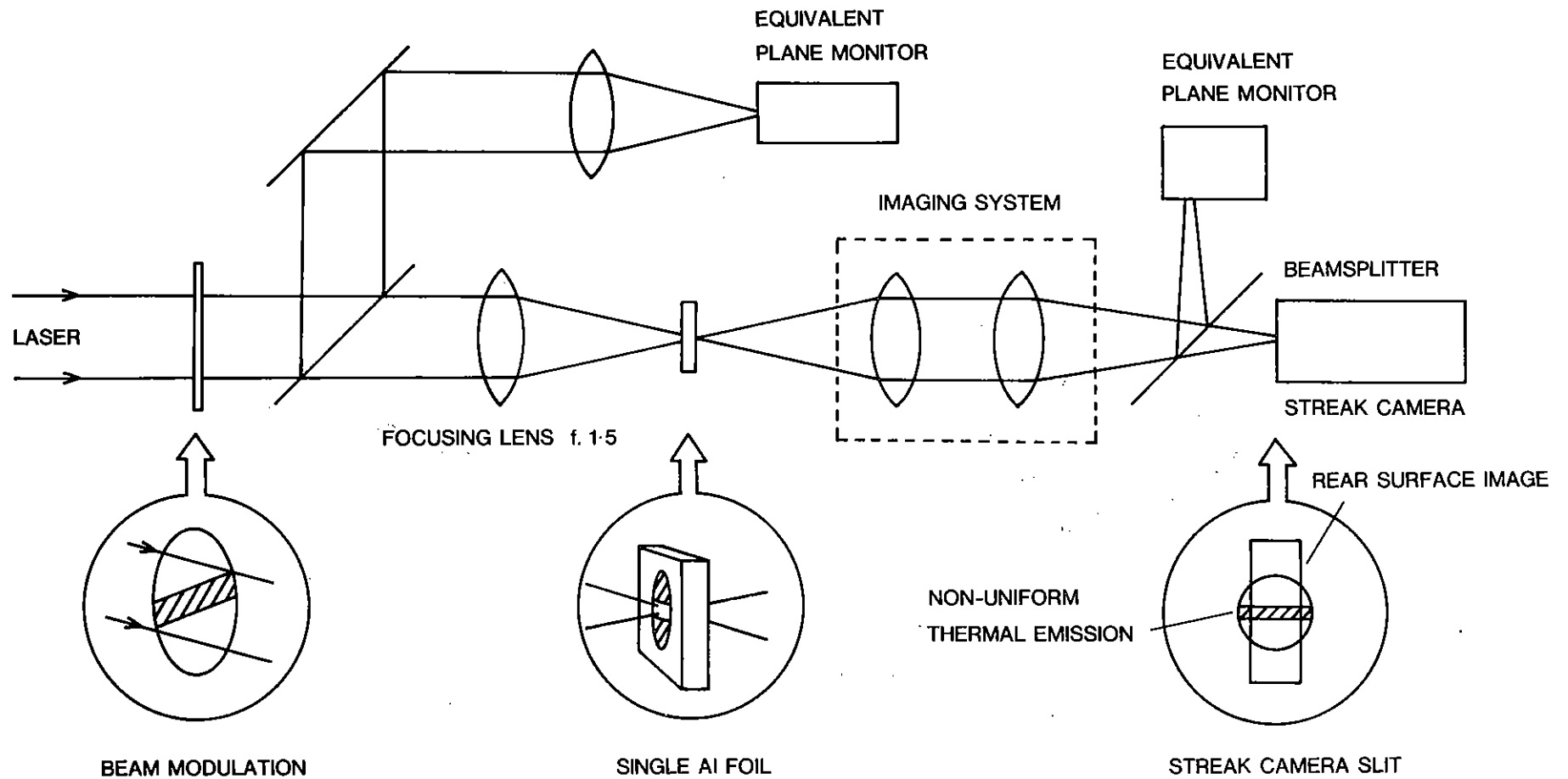


Fig. A2.64 Experimental arrangement for measurement of thermal smoothing using a shock transit technique.

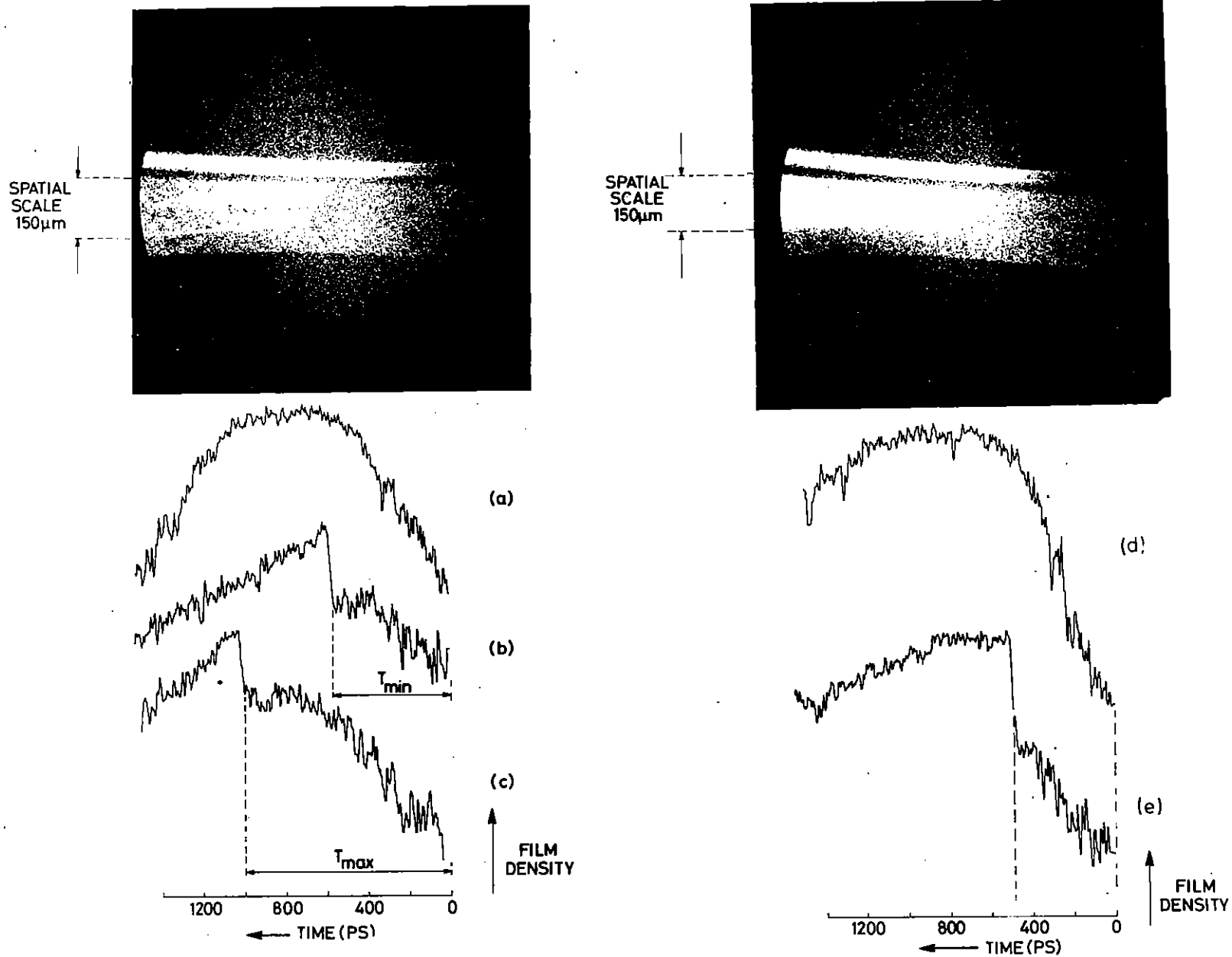


Fig. A2.65 Streak photographs of thermal emission from the rear surface of 25μm Al foils for the case i) modulated laser beam, and ii) unmodulated beam. Together with densitometer tracings of a) and d) laser pulse timing fiducial, b) and c) emergence of shockfront from rear of target at different times corresponding to pressure modulation at ablation front and e) unmodulated shockfront profile.

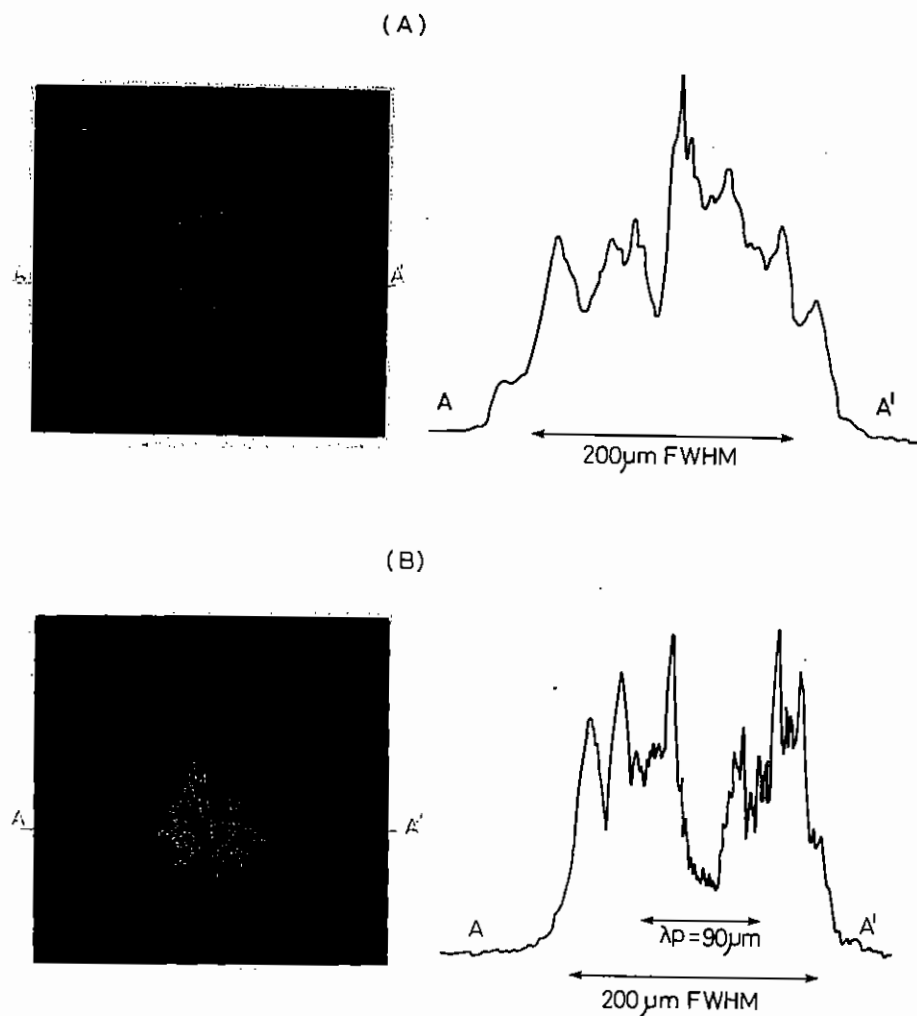


Fig. A2.66 Microdensitometer tracings of equivalent plane photographs for (a) unmodulated beams and (b) modulated beams.

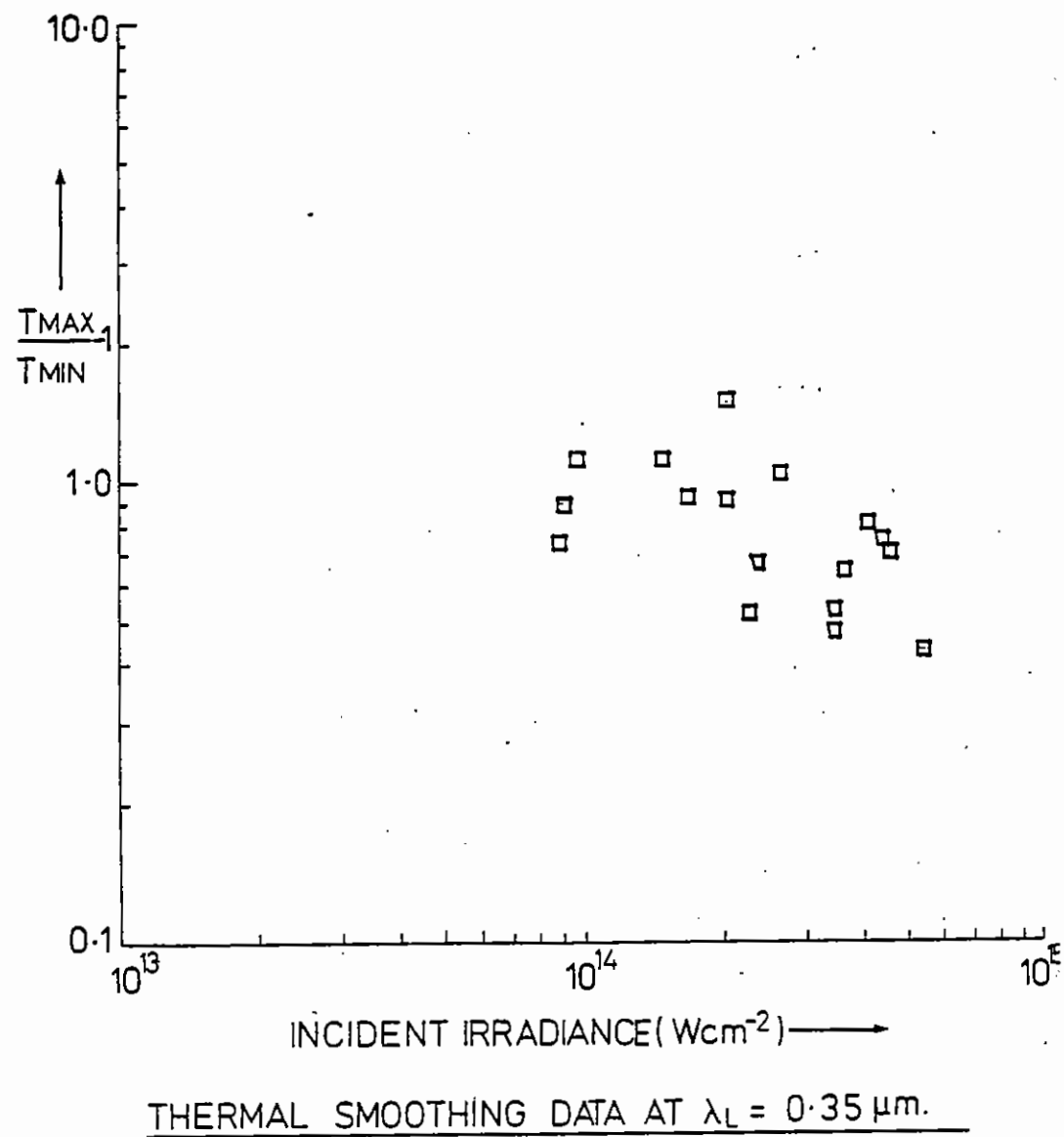


Fig. A2.67 The scaling of $\frac{T_{\text{MAX}}}{T_{\text{MIN}}} - 1$ as a function of incident irradiance at $0.35 \mu\text{m}.$

modulation of $I_{\max}/I_{\min} = 4.5:1$ a zero smoothing limit of 0.75 is obtained, this is close to the observed low irradiance limit.

Only a twofold degree of smoothing has been observed, it seems unlikely that the blue irradiance could be increased significantly taking into account the minimum size of focal spot ($\approx 150\mu\text{m}$) for shock planarity considerations A2.25. Difficulties in obtaining a reliable beam modulation pattern using an interference pattern technique meant that beam modulation scalelengths were limited to those obtained by a neutral density technique. Future experiments will utilise a new interference pattern technique to increase the number of modulation pattern scalelengths over the beam and enter the irradiance region where significant smoothing could be observed. Also a new technique of multilayer target design will enable the depth of shockfront modulation to be observed at different times during the laser pulse.

Filamentation or jet Structures have been observed by a number of workers (A2.26), Figure A2.68 shows a schematic diagram of the pinhole camera arrangement together with a typical photograph and a microdensitometer tracing. The filamentary structures can be seen to extend outside the focal spot region and have a scalelength of $20\mu\text{m}$. These are believed to be the first observations of these structures in this geometry.

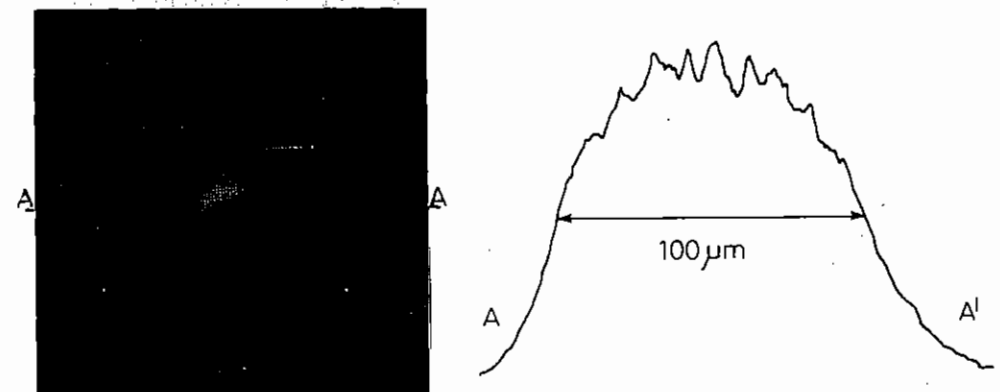
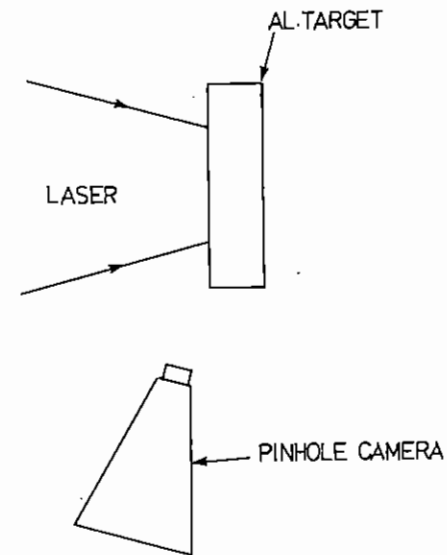


Fig. A2.68 Pinhole camera alignment and photograph together with microdensitometer tracing shown filamentary structures.

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A3 DENSE PLASMA AND RADIATION PHYSICS

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Editor S J Rose

A3.1 Introduction

High-power lasers offer a unique ability to produce in the laboratory material under extreme conditions of temperature and density. Part of the remit of the new group "Dense Plasmas and Radiation Physics" is to study those (non-ideal) plasmas which by virtue of their high density have interparticle interaction energies greater than their kinetic energy and cannot therefore be described by the Debye-Huckel theory which successfully describes more dilute, ideal plasmas. In the last year experiments have been conducted (A3.2) which have demonstrated the ability to produce non-ideal plasmas in planar geometry using the method of colliding shocks, to diagnose the density using a radiographic technique and to record the X-ray absorption spectrum. In addition single-shock propagation has been used to investigate the equation of state of dense plasmas (A3.3).

The other part of the work of the group involves the study of the radiation physics of laser-produced plasmas and work on the identification of satellite lines (A3.4) is reported.

A3.2.1 Colliding shock experiments

J D Hares, D K Bradley, A J Rankin (IC) and S J Rose (RAL)

The aim of these experiments was to compress material to high densities by the collision of two opposed shocks driven by two-sided $0.5\mu\text{m}$ irradiation of a planar target. Planar geometry is particularly amenable to diagnostic probing allowing both streaked radiography and streaked spectroscopic absorption measurements to be carried out.

The radiographic experimental arrangement is shown in Figure A3.1. The targets are self-backlighting by means of the Bismuth strip A. The compression is followed by means of the X-ray shadow of the two buried Bismuth tracer layers B and C. The focal spot was placed on the cross formed by the strips, the rest of the target being transparent. The

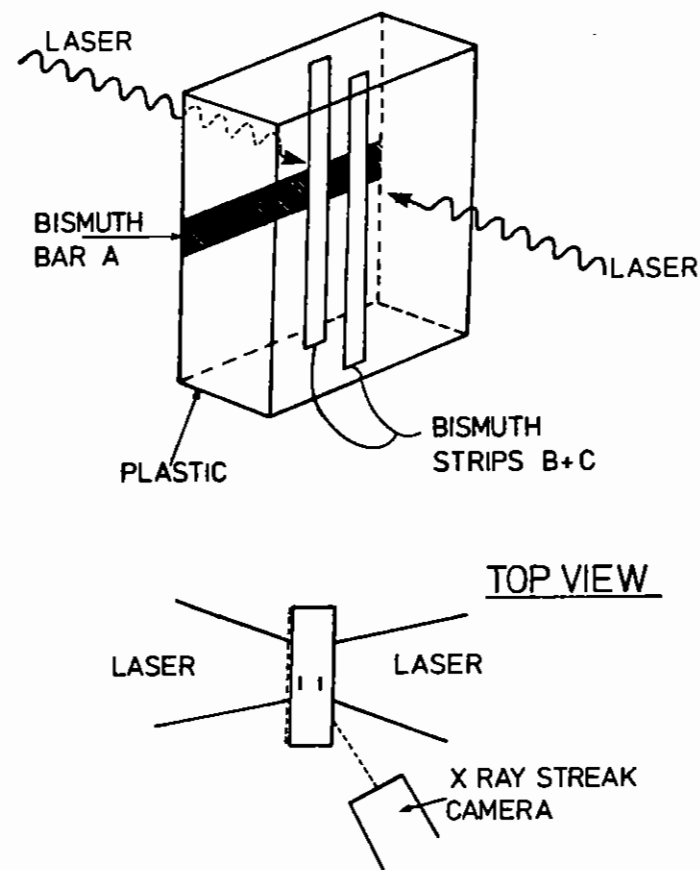


Fig. A3.1

Schematic diagram of the radiographic part of the colliding shock experiment.

angle of view θ , is a tradeoff between target absorption (lower for larger θ) and resolution (better for smaller θ). The angle chosen gave a backlighting cut-off energy of 2.5keV and a spatial resolution (with a $4\mu\text{m}$ pinhole) of approximately $8\mu\text{m}$. Figure A3.2 shows a streaked radiograph from a shot with incident irradiance of $6 \times 10^{13} \text{ W cm}^{-2}$, showing a compression of ≥ 6 , the lower limit being set by the resolution. A simulation of this experiment (using the code MEDUSA) predicts a maximum compression of approximately X 6. The predicted compression from a single shock at this irradiance is approximately X 3 implying a density enhancement of a factor of two by using colliding shocks. The maximum temperature achieved in the simulation (averaged over the central $30\mu\text{m}$ of material) is approximately 11eV.

The major part of the experiment involved K-edge absorption spectroscopy of the plasmas produced using the colliding-shock technique. The K absorption edge of Chlorine was chosen for study, being at a suitable wavelength for both the backlighter and the detector. Chlorine was easily incorporated into the targets by replacing a central plastic layer by chlorinated parylene. The targets were again self backlighting by incorporating a thin layer of Bismuth on one side. The Bismuth was buried under a thin layer of plastic in order to delay emission until maximum compression. In this way the asymmetry produced by the Bismuth layer did not significantly affect the hydrodynamics. Figure A3.3 shows the experimental arrangement together with target details. A Ge crystal spectrometer was used on the front of a miniature reentrant streak-camera. The spectral resolution was approximately 2eV at the Chlorine K-edge. The backlighting continuum intensity was sufficient to record an absorption jump ratio of 5 in the target. The time resolution was approximately 150ps, the incident pulse length being 1.6ns. Figure A3.4 shows the best streak in which a red-shift of the edge is observed. The shift (approximately 10eV) is apparent 300ps after the peak of the laser pulse. MEDUSA predicts the maximum temperature and density to occur 250ps after the peak of the laser pulse although with a focal spot radius of approximately $75\mu\text{m}$ and a shock propagation distance of $30\mu\text{m}$, unlike the simulation, the experiment is not expected to be truly one-dimensional. A comparable red-shift has been seen on several shots.

A3.2

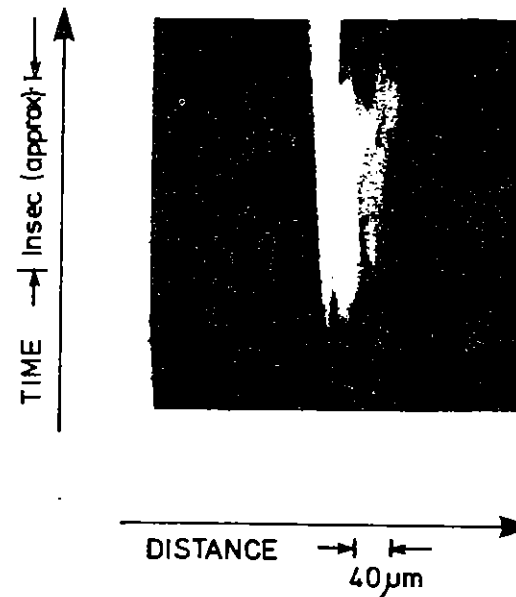


Fig. A3.2

Streaked radiograph from a shot of incident irradiance of $6 \times 10^{13} \text{ W cm}^{-2}$.

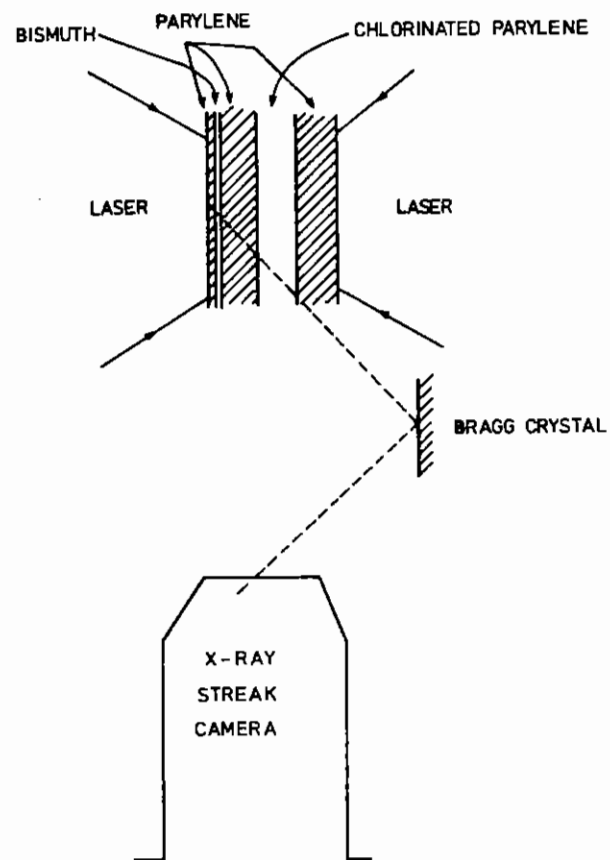


Fig. A3.3 Schematic diagram of the spectroscopic part of the colliding shock experiment.

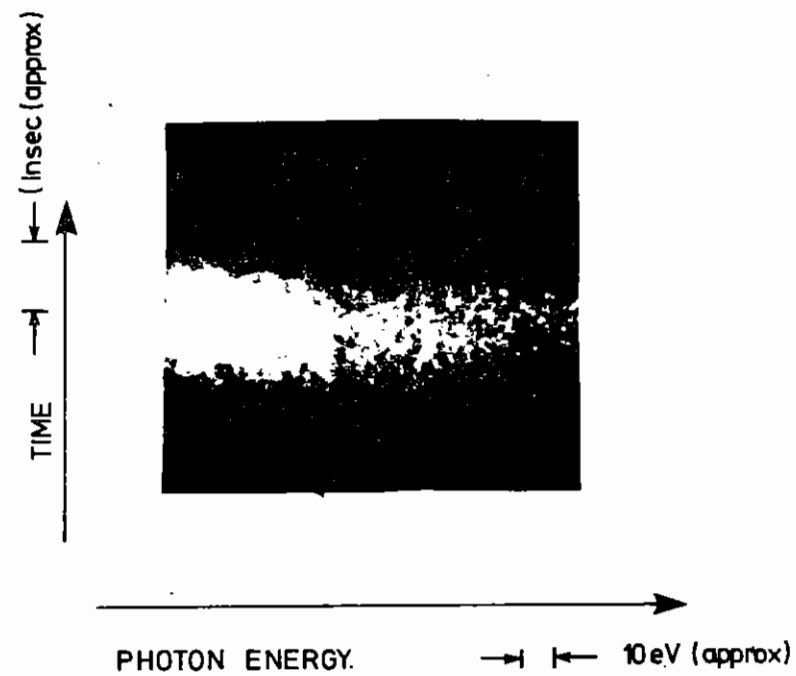


Fig. A3.4 Streaked spectroscopic record showing red shift of chlorine K-edge.

A prediction of the shift in the Chlorine K-edge has to include several effects. Firstly, as the material increases in temperature and density ionisation occurs. This leads to an increase in the K-shell ionisation energy. However opposed to this is the depression of the continuum which goes to lower the ionisation energy. Finally at peak density and temperature the material is strongly degenerate which shifts the onset of photoionisation to higher energy. A simple model has been constructed to include these effects which estimates the ionisation in the material from Thomas-Fermi theory and uses relativistic Hartree-Fock calculations to predict the change in the ionisation energy with increasing ionisation. Without inclusion of the effects of degeneracy the model predicts a shift of -33eV but inclusion of the Fermi energy reduces the shift to -11eV , in good agreement with experiment. The model also predicts that the plasmas produced in these experiments are strongly-coupled with $\Gamma \sim 2$.

The experimental investigation reported here is one of several characterising the behaviour of K-edges in plasmas. Overall two notable observations have been made.

- (i) Temperature increase at solid density results in broadening of the K-edge without shift.
- (ii) Density increase results in a red shift. It is qualitatively reasonable that if the density remains as solid that the edge position should not shift with temperature increase because, although the ionisation would be expected to produce a blue shift, the free electrons are still present in the same volume of space. The broadening results from statistical variations in the microfield and is consistent with a Holtzmark-type distribution. A density increase however results in an imbalance between the competing shift mechanisms of ionisation and continuum lowering. In the experiments performed here, the temperature remains low (approximately 11 eV from simulation) and the broadening is relatively small.

In conclusion, these experiments have demonstrated the ability to perform absorption spectroscopy of non-ideal degenerate plasmas. Simple

hydrodynamic and atomic physics models give results in good agreement with experiment.

A3.2.2 Particle/Shock Velocity Measurement

An important first objective of the colliding shock experiment was to measure the particle and shock velocity of the buried Bismuth strip targets using single-sided irradiation. As in the two-sided irradiation case the trajectories of the Bismuth tracer strips could easily be followed using X-ray streaked shadowgrams.

Figure A3.5 shows a schematic diagram of the strips and their reaction as the shock encounters the first strip and then, after a small delay the second. A typical streak shadowgram shows the high resolution and good clarity obtainable using this technique.

The shock velocity is simply calculated from the time delay between the first strip being hit by the shockfront and the subsequent movement of the second strip as it too reacts as the shock hits it. Knowing the separation of the Bismuth strips gives the average shock velocity between the two strips. The particle or material velocity is easily calculated from the trajectory of the first foil.

Complications arise in the interpretation of these trajectories as pincushion distortion effects in the camera intensifier distort the image paths. A technique will be applied to deconvolute this non-linear effect from digitized densitometer scans of the pictures. However, the uncorrected photographs give typical values of

$$\begin{aligned}\text{shock velocity} &\sim 2-3 \times 10^6 \text{ cm s}^{-1} \\ \text{particle velocity} &\sim 5 \times 10^5 \text{ cm s}^{-1}\end{aligned}$$

Shock Hugoniot give the drive pressure to be approximately 18 Mbar .

These excellent results give the first simultaneous shock and particle velocity measurements. A future experiment would use materials whose equation of state is better known. This would then allow a comparison to be made between the pressure calculated from the particle velocity and that

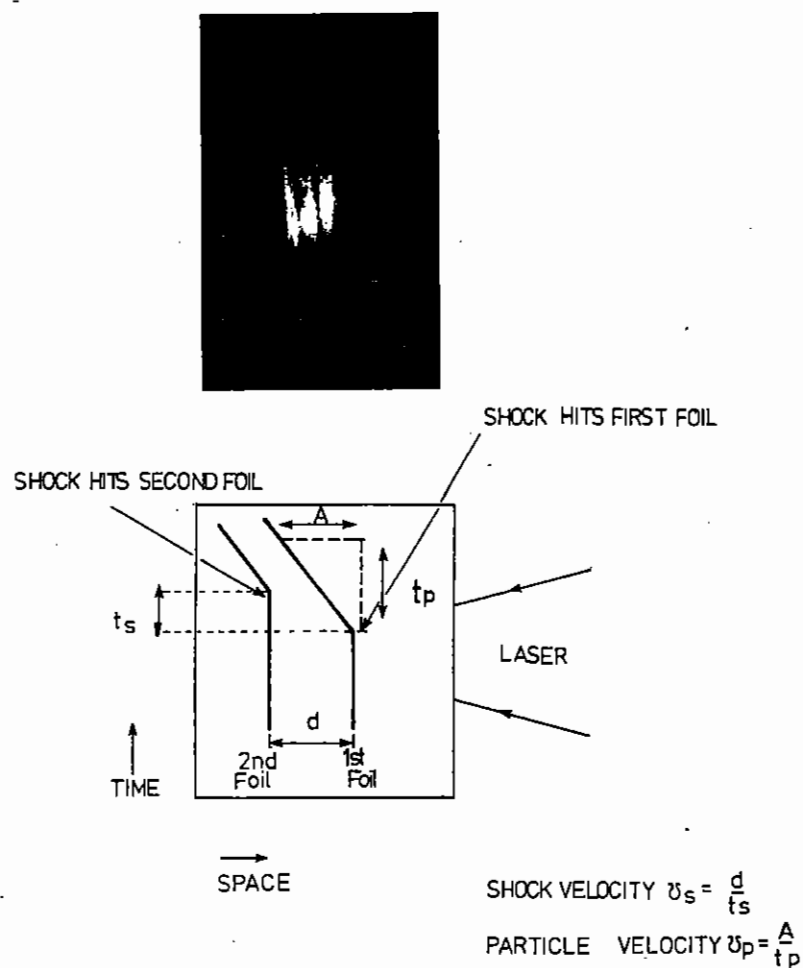


Fig. A3.5

Typical streaked spectrograph of radiographic measurement under single-sided irradiation.

from shock pressure inferred from shock velocity. Some of the discrepancies between ablation pressure (measured from ablation rate) and shock pressure (measured from optical streaks of stepped targets) may then be resolved.

A3.3 Shock propagation experiments at $0.53\mu\text{m}$

A J Rankin (IC)

A3.3.1 Introduction

It is widely recognised that an important diagnostic approach to measure parameters in laser plasma experiments such as ρ , T and P can be the detection of laser generated shockwaves as they transit the target and emerge through the back. The arrival of the shockfront at the rear surface is signalled by prompt heating of the material and can be detected easily by the use of an optical streak camera to monitor the continuum emission in the visible part of the spectrum. A number of experiments have been performed in many labs (A3.1, A3.2, A3.3) to measure the shockwave velocity and through the E.O.S. infer pressure in the shocked material. A few anomalous results were obtained at AWRE Aldermaston (A3.4) at $1.06\mu\text{m}$ but until now no other measurements have been obtained in this country. The experiments described in this section signal the first attempt to measure shockwave velocity at this laboratory and constitute one of the most comprehensive studies of shockwave propagation through both stepped and planar foil targets consisting of both single and multilayer materials.

A3.3.2 Experimental

A single beam of the VULCAN Nd glass laser system was frequency doubled to irradiate stepped and planar foils with $0.53\mu\text{m}$ radiation of 1ns FWHM pulse duration. Figure A3.6 shows a diagram of the experimental arrangement. Stepped Aluminium targets, with typically a $12\mu\text{m}$ base and a $6\mu\text{m}$ step were aligned using a CW beam axis reference together with a rear surface illumination system to locate the step onto the $100\mu\text{m}$ wide streak camera slit, orientated transversally to the step plane. The imaging system consists of 25 mm microscope objective relayed through light shields to a

SHOCK VELOCITY MEASUREMENTS USING STEPPED FOIL TARGETS

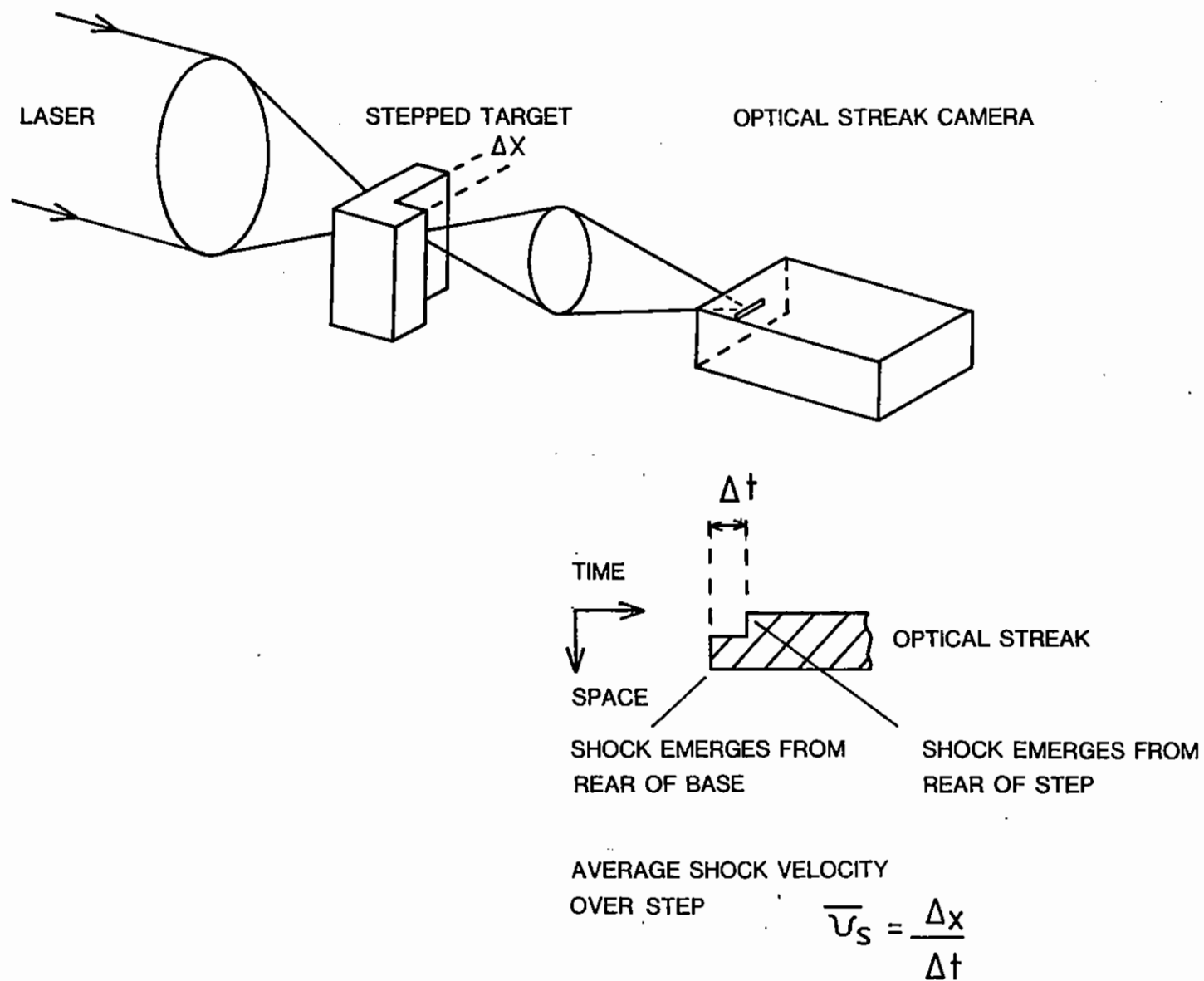


Fig. A3.6

Schematic diagram of the experimental arrangement to measure the average shock velocity over an aluminium stepped target.

38cm achromatic camera lens giving an overall image magnification of 25 onto the Hadlands streak camera, fitted with an S20 photocathode. A small fraction of the $0.53 \mu\text{m}$ radiation was relayed onto the camera slit via a 1m length of single-mode plastic 1mm fibre-optic to give a timing fiducial mark. The planar targets consisted of layers of transparent material with an initial thin Bismuth coating on the front surface. (This design was governed by the mainstream experimental requirements.) A few planar Aluminium targets were also irradiated for comparison with the stepped target shots. Focal spots were of the order of $150\mu\text{m}$ to avoid significant shockwave non planar effects (A3.5). The streak photographs for a typical unstepped and stepped Aluminium target are shown in Figure A3.7. There is some evidence of beam nonuniformity in the stepped target emission which was also measured by the equivalent plane photograph for that shot. The average shock velocity over the step has been measured for irradiances of $2 \times 10^{13} - 3 \times 10^{14} \text{ W cm}^{-2}$. Using SESAME E.O.S. data for aluminium (A3.6), a graph of shock pressure as a function of absorbed irradiance has been plotted in Figure A3.8.

Figure A3.9 shows a multilayer planar target together with a typical optical streak photograph. In the image, shock generated thermal emission is detected early in time by the cameras 'through' the following 'transparent' layers of material as the shockfront transits the initial Bismuth layer. A period then follows of zero or low-level emission beyond the sensitivity range of the camera when the shockfront may be transiting the mylar layer. The low-level emission may be due to radiation heating of the mylar by the Bismuth, increasing its opacity and preventing the emission being detected. The shock then crosses the Mylar KCl boundary where the level of emission changes abruptly. This level is monitored until a further abrupt change signals the shockfront crossing the KCl-CH boundary. Emission then continues until the shock breaks out the back of the target into vacuum. In this particular case the focal spot is very small ($\sim 50\mu\text{m}$) and gross 2D non-planar evolution is evident. Other shots with targets of different

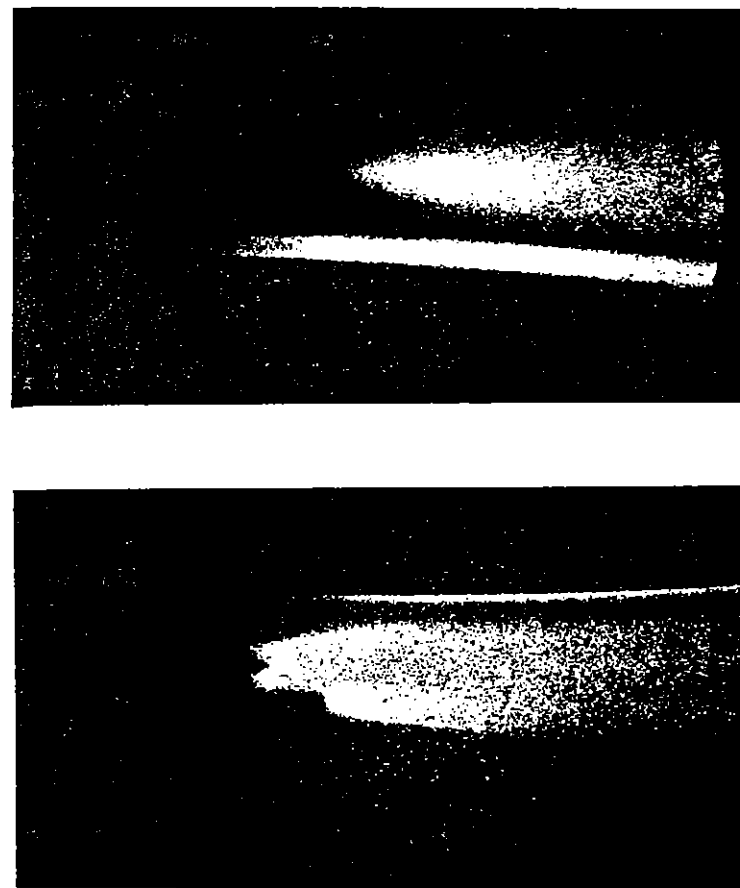


Fig. A3.7

Optical streak photographs of an unstepped and a stepped aluminium target ($12\mu\text{m}$ and $6\mu\text{m}$) at an incident irradiance of approximately $10^{14} \text{ W cm}^{-2}$.

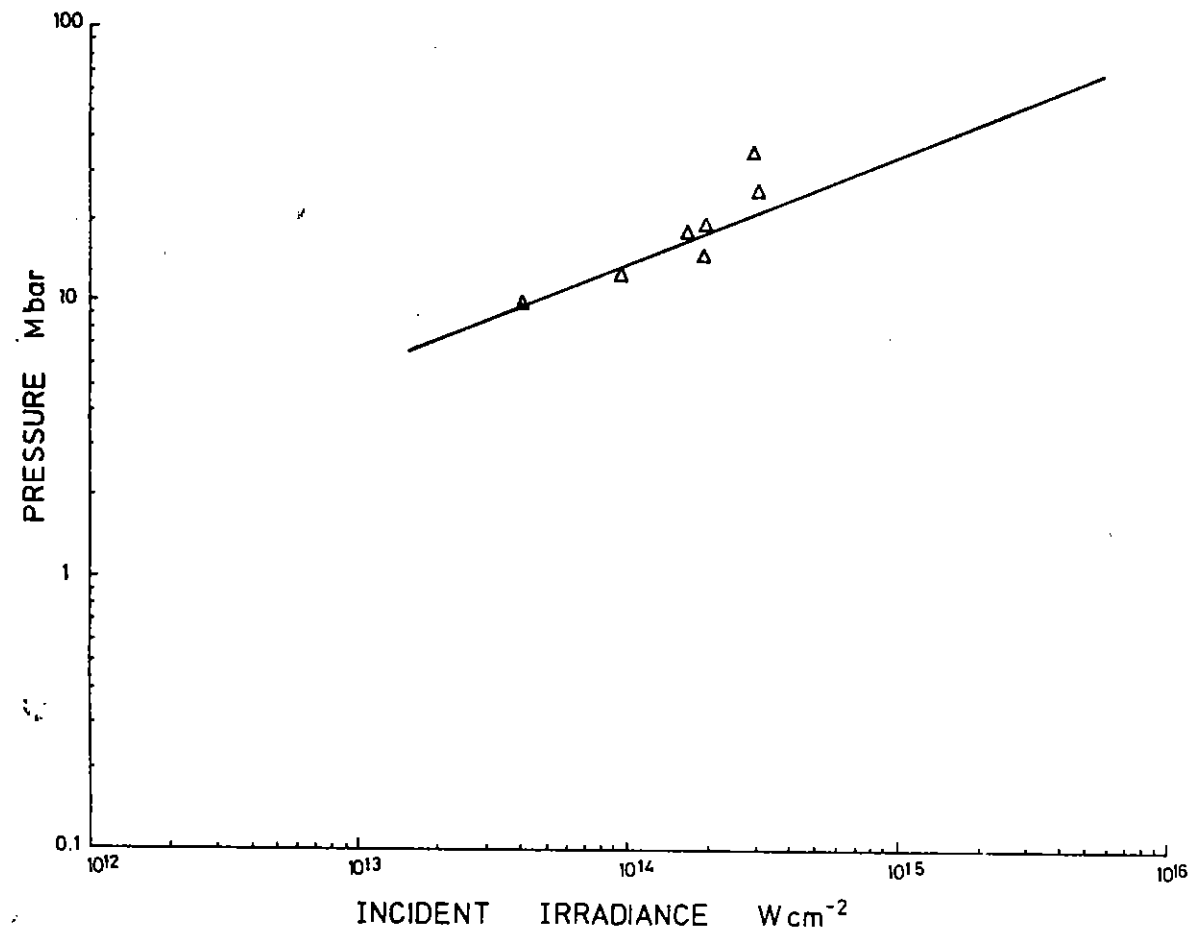
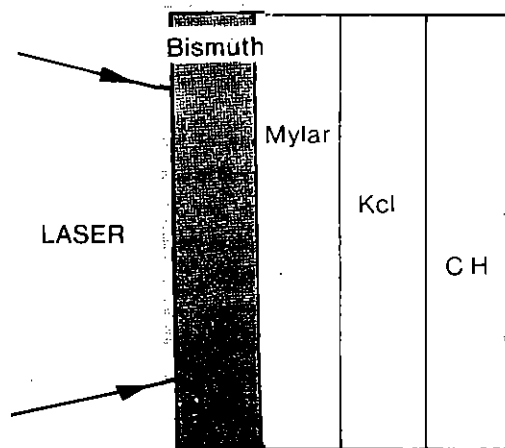


Fig. A3.8 A graph of inferred shock pressure (Mbar) as a function of absorbed irradiance for 0.53 μ m irradiation.



dimensions show larger transit times which scale sensibly. More experiments are required in a mainstream rather than a parasitic capacity. Comparison of these results with simulations performed recently (A3.7) on shock propagation across target interfaces may yield information on the compression state of the different target layers.

The stepped target experiments have produced good stepped streaks which have allowed accurate measurements of the average shock-velocity in Aluminium targets. The success of the experiments to predict pressure values which are comparable to previous studies is extremely useful.

The multilayer target results demonstrate a new technique to observe the spatial shape and position of the shockfront through the duration of the laser pulse. This may be a technique to measure average shock-velocity in a planar target avoiding the complication of stepped targets and their inherent problems.

The co-operation of D K Bradley and R W Eason is acknowledged in the acquisition of this data during the mainstream experimental time. Also the assistance of the AWRE target preparation group in manufacturing the aluminium targets for mounting here at RAL is gratefully acknowledged.

A3.4 One-Photon two-electron dielectronic satellite transitions in the X-ray spectra of laser-produced plasmas

J G Lunney (Trinity College, Dublin)

X-ray spectroscopy is now widely used to measure the temperature, density and state of ionisation of plasmas produced by the laser ablation of solid surfaces and the laser implosion of microballoon targets. This work has recently been reviewed by De Michelis and Mattioli (A3.8). Dielectronic satellite lines close to the hydrogen-like and helium-like resonance lines are produced by the radiative decay of doubly-excited autoionising levels. The intensities of these satellite lines relative to the resonance lines depend on the plasma temperature (A3.9, A3.10). At high electron density, electron collisions redistribute the populations of the autoionising levels

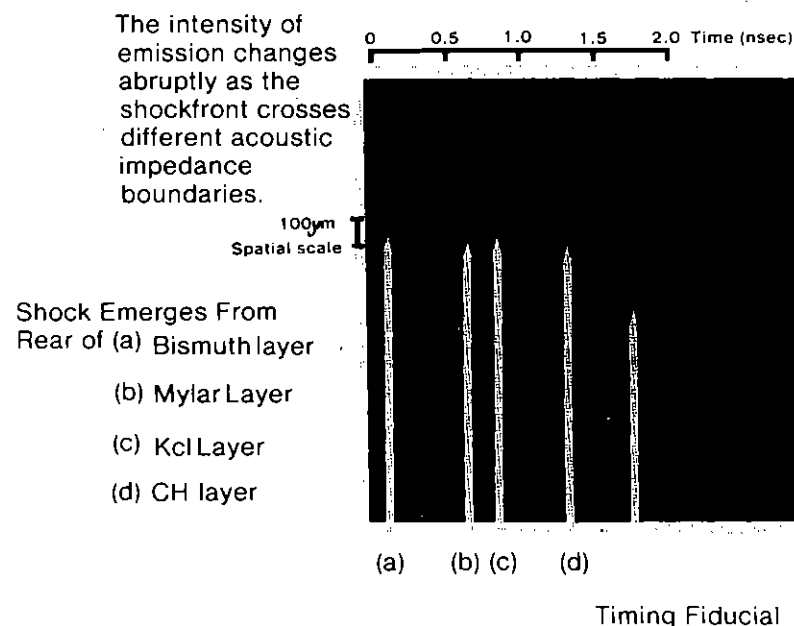


Fig. A3.9 Schematic diagram of a multi-layer target together with an optical streak photograph.

and make the relative intensities of the dielectronic satellite lines density dependent (A3.11 to A3.14). Thus, dielectronic satellite lines, which occupy a small spectral range, can be readily applied to the measurement of temperature and density in laser-produced and other plasmas. This paper reports the observation of the one-photon two-electron dielectronic satellite lines Al XII $1s2p\ ^1P_1 - 2s^2\ ^1S_0$ and Al XI $1s^2 2p\ ^2P_{1/2,3/2} - 1s2s^2\ ^2S_{1/2}$ in the X-ray spectra of laser-produced plasmas. The lines $1s^2 2p\ ^2P_{1/2,3/2} - 1s2s^2\ ^2S_{1/2}$ are labelled (o,p) by Gabriel (A3.9). The $1s2p\ ^1P_1 - 2s^2\ ^1S_0$ line was also observed for magnesium and silicon. Previously, the (o,p) satellite line has been observed for lithium-like boron in the spectrum of a laser-produced plasma (A3.15).

The plasma was produced using the frequency-doubled output of one beam of the Nd-glass laser system. The second harmonic beam had an energy of 21J and a pulse length of 1 ns FWHM. The beam was focused at nearly normal incidence on a plane aluminium target to a focal spot diameter of 80 μ m giving a peak irradiance of about 4×10^{14} Wcm $^{-2}$.

The X-ray spectra were recorded using a miniature space-resolving X-ray spectrograph (A3.16) which viewed the plasma in a direction making an angle of approximately 10° with the plane of the target surface. An ammonium di-hydrogen phosphate crystal (ADP, $2d=10.648\text{\AA}$) was used, giving a spectral range of 4.8-9.0 $\text{\AA}$.

The spectra were recorded with a single laser shot. A 25 μ m slit lying in the plane of dispersion was used to provide spatial resolution normal to the target surface, and so separate the emission of the ablation and recombining phases of the plasma. Analysis of the spectrum shows that in the ablation plasma the electron temperature is 440 ± 20 eV and the electron density is $(1 \pm 0.2) \times 10^{22}$ cm $^{-3}$ (A3.16).

The Al XIII L_α and its associated satellites are shown in Figure A3.10. The one-photon two-electron satellite Al XIII $1s2p\ ^1P_1 - 2s^2\ ^1S_0$ line is observed at 7.368 $\text{\AA}$. Figure A3.11 shows the Al XII $1s^2\ ^1S_0 - 1s2p\ ^1P_1$ line with its associated satellites including the blended (o,p) line at 8.078 $\text{\AA}$. In Table A3.1, measured wavelengths of these one-photon

two-electron satellites are compared with the calculations of Vainshtein and Safronova (A3.17). Taking into account the relative intensities of the two components of the (o,p) lines, the theoretical wavelength for the centroid of the blended line is 8.0779 $\text{\AA}$. In these wavelength measurements, the other more intense satellites were used for wavelength calibration. The spectral resolution was limited by the radial dimension of the plasma, and could be improved by using a smaller laser focal spot or by increasing the distance between the plasma and the film. Better wavelength calibration could be obtained by doping the target with a suitable lower Z element concentration to give optically thin He-like and H-like resonance lines. The $1s2p\ ^1P_1 - 2s^2\ ^1S_0$ line was also observed in spectra recorded from magnesium and silicon laser-produced plasmas and the measured wavelengths are shown in Table A3.1.

Table A3.1 Observed and calculated wavelengths

Transition	Wavelength($\text{\AA}$)	
	Calculated	Observed
Mg XI $1s2p\ ^1P_1 - 2s^2\ ^1S_0$	8.6684	8.668 ± 0.003
AlXII $1s2p\ ^1P_1 - 2s^2\ ^1S_0$	7.3690	7.368 ± 0.003
AlXI $1s^2 2p\ ^2P_{1/2} - 1s2s^2\ ^2S_{1/2}(p)$	8.0754	8.076 ± 0.003
AlXI $1s^2 2p\ ^2P_{3/2} - 1s2s^2\ ^2S_{1/2}(o)$	8.0792	
SiXIII $1s2p\ ^1P_1 - 2s^2\ ^1S_0$	6.3409	6.340 ± 0.003

At sufficiently low electron density to make collisional redistribution between autoionising levels unimportant, the populations of these levels are mainly determined by dielectronic recombination. The intensity I_{ij} of the satellite line from level i to level j due to dielectronic

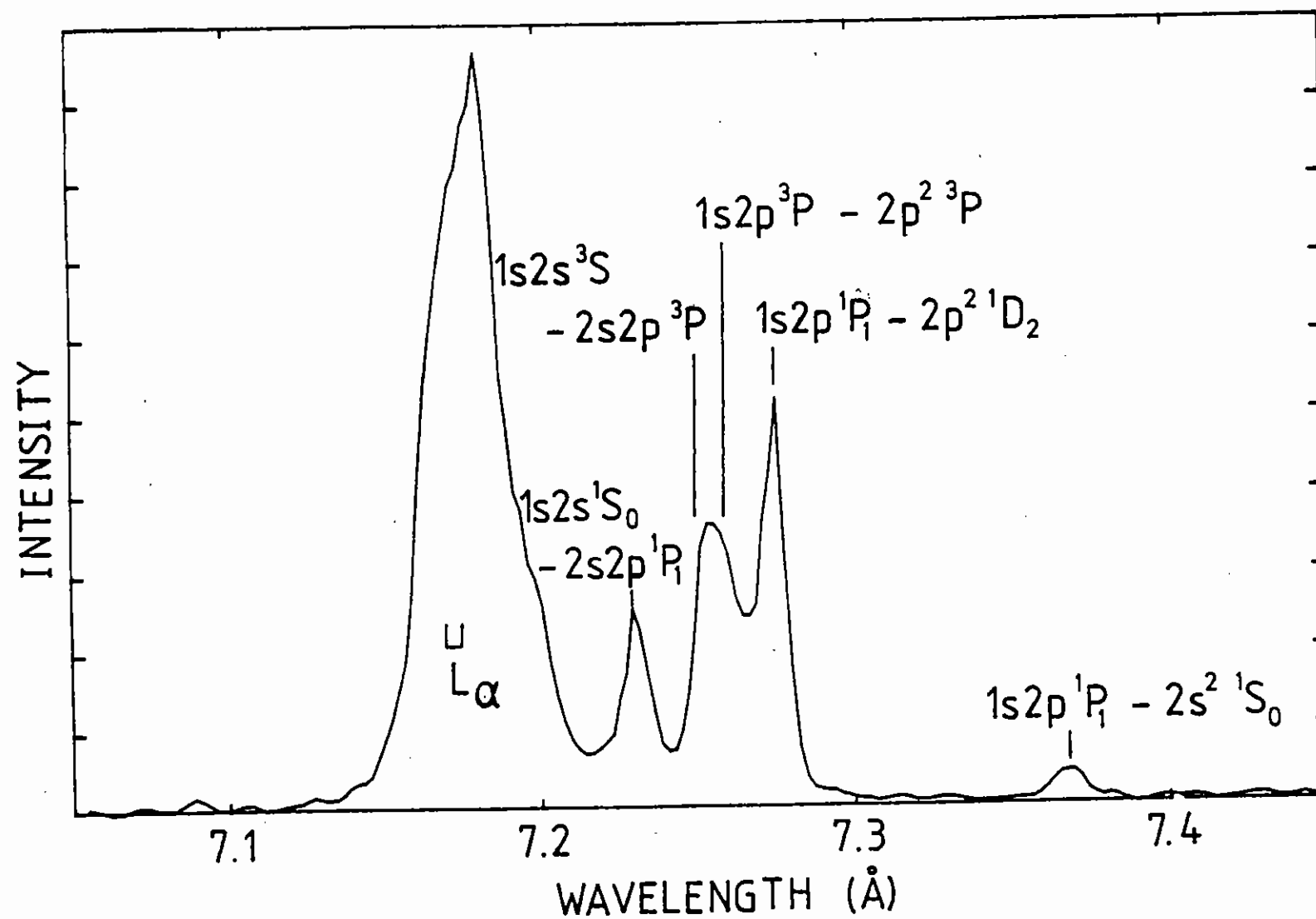


Fig. A3.10 The AlXII $L\alpha$ line and its associated satellites including the AlXII $1s2p^1P_1 - 2s^2^1S_0$ line at 7.368 Å.

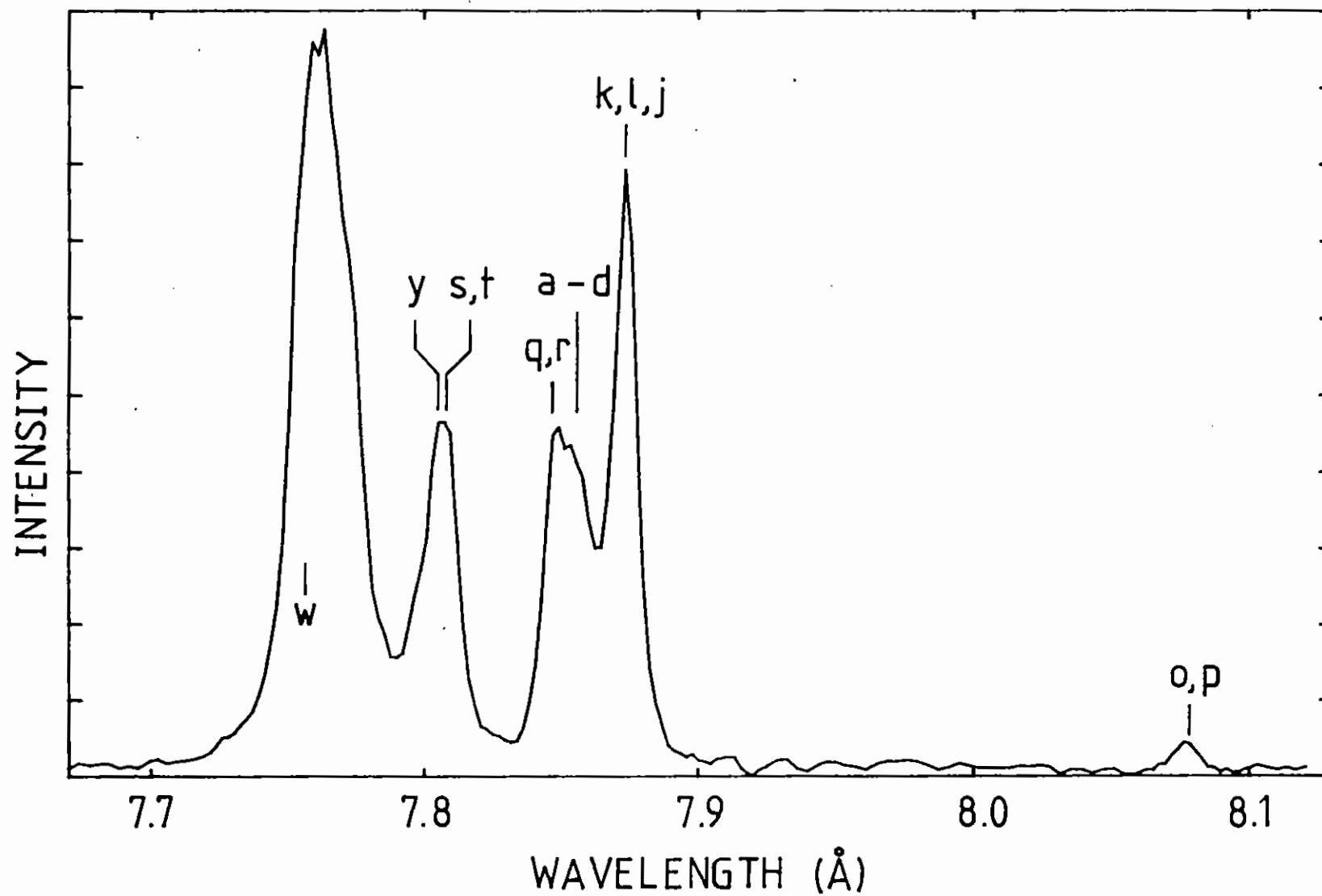


Fig. A3.11

The AlXII $1s^2\ ^1S_o - 1s2p\ ^1P_1$ line (w) and its associated satellites including the (o,p) satellite at 8.076 Å. The satellite lines are denoted using the key letter description of Gabriel (A3.10).

recombination is proportional to the factor q_{ij} given by

$$q_{ij} = \frac{g_i |A_{ij}|}{(|A_i| + \sum_k A_{ik})}$$

where $|A_i|$ is the autoionisation rate and A_{ij} is the radiative decay rate (A3.9). The relative intensities of lines derived from levels with a high autoionisation rate are not significantly changed by collisions between autoionising levels at high electron density (A3.12). For the satellites of the aluminium L_α line Vainshtein and Safronova (A3.17) give $q(1s2p^1P_1 - 2s^2^1S_0)/q(1s2s^1S_0 - 2s2p^1P_1) = 0.14$ and $q(1s2p^1P_1 - 2s^2^1S_0)/q(1s2p^1P_1 - 2p^2^1D_2) = 0.044$ while the observed line intensity ratios are 0.16 and 0.08 respectively. The high relative intensity of the $1s2p^1P_1 - 2s^2^1S_0$ line in the experimental spectra is probably due to the self-absorption of the other stronger satellite lines. This is corroborated by the fact that $q(1s2p^1P_1 - 2p^2^1D_2)/q(1s2s^1S_0 - 2s2p^1P_1) = 3.2$ while the experimental intensity ratio is 2.0. The opacity of the $1s2p^1P_1 - 2p^2^1D_2$ line is about 3.3 times that of the $1s2s^1S_0 - 2s2p^1P_1$ line. Also, the $1s2p^1P_1 - 2s^2^1S_0$ line was not observed in spectra recorded from targets of 0.1 μm of aluminium on Mylar. Thus, the relative intensities of these singlet dielectronic satellite lines can give information about the density $\times$ length product of the plasma being observed.

For the satellites of the AlXII $1s^2^1S_0 - 1s2p^1P_1$ line $q(o,p)/q(k,i,j) = 0.031$ while the experimental intensity ratio is 0.06. Again, this discrepancy is due to self-absorption of the stronger lines. Also, the (o,p) satellite was not observed in spectra from thin layer targets made with 0.1 μm of aluminium on Mylar.

In conclusion, the one-photon two-electron dielectronic satellite lines have been observed in magnesium, aluminium and silicon laser-produced plasmas. The measured wavelengths agree with theoretical predictions to within our wavelength resolution. The intensities of these two-electron satellite lines relative to other one-electron satellites show that there is significant self-absorption of the one-electron satellites.

Section A3

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A4 XUV LASERS AND APPLICATIONS OF LASER PRODUCED PLASMA

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A4 XUV LASERS AND APPLICATIONS OF LASER-PRODUCED PLASMAS

A4.1 Introduction

G J Pert (Hull)

In the previous year the group has investigated two approaches to XUV laser action, namely recombination in the carbon fibre scheme, and photo-excitation using the seeded foil approach.

The carbon fibre experiments have been successful in confirming the model of the interaction observed in computer simulation, and it is clear that the underlying physics of this scheme is now well understood. However, it has not proved possible in these experiments to consistently measure gain. It is now clear that there exist two modes of operation: one at low absorbed energy is of high gain, but of a very narrow window and was observed in the early experiments; the other of lower gain has a larger window and is more suited to general use. At present the absence of sufficiently good line focusing prevents the construction of a working laser in either mode. Only limited further work on this scheme is therefore likely to be carried out until this limitation can be improved.

In the coming year it is therefore probable that the major effort will concentrate on the photo-excitation scheme. Preliminary experiments to observe pumping were somewhat inconclusive, but yielded valuable experience in construction, and handling of the foils. In the coming year studies should enable the detailed physics of the pumping to be explored and understood.

The second major area of activity of the group is the generation of incoherent X-rays as a source for other work such as X-ray microscopy and X-ray absorption spectroscopy. Notable contributions to this field have been made by the group in the past and there remains an important area of group activity. In the past year a first attempt was made to use a laser source for an X-ray microscope instead of the more conventional synchrotron. Further experiments in the coming year should enable the practicability of this scheme to be evaluated.

There is little doubt that the high powered laser as a generator of incoherent X-rays remains poorly exploited. In comparison with synchrotron sources it offers a high flux, short pulse duration source of as yet poorly defined spectral quality. However, until it has been systematically surveyed its full potential, and utilisation remain unclear. The identification and gradual accomplishment of this goal must remain one of the long term objectives of the group.

A4.2 CVI Experiments

M J Lamb, C L S Lewis, R Corbett (QUB), L D Shorrocks, G J Pert (Hull), M H Key, C J Hooker, R W Eason (RAL) and Fan Pinzhong (Shanghai)

A4.2.1 Introduction

The use of thin carbon fibres irradiated by cylindrically focused Nd:glass laser radiation, and its harmonic as a working medium for laser action in the XUV has been investigated both theoretically (A4.1) and experimentally (A4.2, A4.3) for some years now. Two sets of experiments have been carried out, one using single beam illumination and 1.06 μm radiation at relatively low energy (A4.2) and the second with four beams and 0.53 μm wavelength, at higher energy (A4.3). These two experiments gave markedly different results, high gain being observed (A4.2) at low energy and low gain (A4.3) at high energy. The interpretation (A4.2, A4.3) of these results in terms of the effective mass of fibre clarified this apparent discrepancy, but in the absence of experimental data on the magnitude of the burn, the detail remained obscure. The recent experimental measurements of the effective (burn) mass in both configurations allow us to confirm the predicted behaviour, and to check the code values.

The high gain values observed at low energy (A4.2) are necessary to construct a working laser, however, the energy window over which they occur allows a very narrow tolerance in the input conditions, probably too small for regular application. This window can be significantly widened by a suitable pump configuration which matches the burn and gain condition. The restricted fibre burn thus has two extremely important consequences.

- (a) High gain can be achieved with large diameter fibres.
- (b) Wide gain windows can be found.

The first of these avoids the problems of handling very thin (2 μm diameter) fibres and allows the use of 4 - 5 μm ones, which are relatively straightforward. The second may overcome some of the problems of laser shot-to-shot variation, and the design and construction of high quality cylindrical optics giving a line focus of 1-2 cm in length.

A4.2.2 Scaling Laws and Operating Conditions

It is clear from the above remarks that there is only a restricted region of parameter space in which high gain can be achieved. If the plasma is too cold severe recombination to the ground state gives rise to strong reabsorption of the Lyman α line; on the other hand if too hot recombination population of the upper states is weak. Thus we would expect that for constant plasma dimensions there is an optimum input energy for peak gain. Furthermore since the opacity limitation is weakened, both directly and due to larger velocity gradients, by smaller transverse dimensions, the gain may be increased by using smaller initial fibre mass per unit length, M . This behaviour has been confirmed by computer modelling of uniformly heated plasma cylinders (A4.1), where it was found that the maximum gain scaled roughly as

$$G \approx \text{const } M^{-17/22} \quad (1)$$

and the energy per unit length necessary to achieve optimum gain scaled as

$$E \approx \text{const } M^{17/11} \quad (2)$$

for a constant heating pulse profile. At very low masses departures from these scalings occur when the expansion is too rapid to permit complete ionisation to be achieved: typically this limits the mass to that of a fibre of diameter greater than 1.5 μm .

Practical realisation of such a scheme requires a heating pulse for the carbon of about 20 J/cm in a time of 100 ps for a mass of 3 μm equivalent diameter. Such energy densities are readily achieved by laser irradiation using a mode-locked Nd:glass system focused on to a line. The limited gain window in pump energy places limitations on the uniformity of the line focus. For a typical case the half gain width in energy allows a variation of $\pm 25\%$ about the optimum. This is typical of the tolerance allowable in the focus for completely burnt fibre. In practice such uniformity is difficult to achieve with the simple two lens arrangements (A4.2, A4.3) which have been used up to now. In consequence experiments have been limited to fibre lengths of less than 3 mm.

During the heating phase absorption of the laser radiation occurs at a density much below that of the solid. The fibre is thus surrounded by a tenuous, hot plasma blanket, which if the initial diameter is sufficiently small completely envelopes the remaining solid. Heat is transferred by thermal conduction radially into the solid thereby forming the hot plasma by ablation: due to the strong non-linear temperature dependence of the thermal conductivity there is a sharp interface between the hot and cold components of the fibre (A4.2). Since the radius of the hot plasma is much larger than that of the initial fibre the hot plasma closely approximates that of a uniformly heated fibre, of the same mass. Following an incomplete burn we may identify two components namely the cold residual core, and the hot plasma; in the latter we may expect an expansion phase having inversion properties similar to those of equivalent uniformly heated fibre of the same mass and energy. Since the heat flow is strongly non-linear the temperature is nearly constant at a value sufficient to ensure complete ionisation in the hot plasma, independent of the fibre diameter, and laser pulse energy, and the ablated hot mass therefore scales roughly proportionally to the absorbed energy

$$E \approx \text{const } M \quad (3a)$$

In practice the temperature must increase slightly as the mass increases, so that (3a) represents a lower bound power law relation.

If the pulse length is sufficiently long the thermal conduction dominated ablation zone between the absorption surface and solid fibre is steady state. Since density of the solid is much greater than that at absorption, and the radius of the initial fibre much smaller than that of the absorption surface, the fibre behaves as a point mass source at the origin. The radius of the absorption surface, R , at which the bulk of the laser radiation is absorbed in some arbitrary manner must be proportional to the thermal conduction length (A4.4) at the point. Hence

$$R = \text{const } T^2/N \quad (3b)$$

where N and T are the density and temperature at the ablation surface respectively. The mass ablation rate per unit length

$$m = \text{const } NT^{1/2}R = \text{const } T^{5/2} \quad (4)$$

and the energy outflow rate, which must balance the laser energy absorption rate

$$e = \text{const } mT = \text{const } T^{7/2} \quad (5)$$

Hence for a constant laser pulse length

$$E = \text{const } M^{7/5} \quad (6)$$

When the steady state condition is not satisfied the ablation zone will be shorter than in steady flow, and the heat transfer to the solid more rapid. For a given temperature the mass ablation rate is increased, or alternatively a lower temperature required to achieve a given mass rate. Under these circumstances the constant temperature condition is more closely obeyed. Thus in general we may write

$$m = \text{const } T^n \quad ; \quad e = \text{const } T^{n+1} \quad (7)$$

or

$$E = \text{const } M^{(n+1)/n} \quad (8)$$

where $1 < n < 1.4$, the lower bound being appropriate for short pulses, and the upper for long pulses, with respect to the characteristic time R/c , c being the sound speed at the ablation surface. The proceeding argument does not depend on the specific absorption process. In contrast the actual conditions at the ablation surface are determined by the details of the absorption. Thus if, as is usually the case, inverse bremsstrahlung dominates, a self-regulating condition $aR \sim 1$ yields a relation $NT^{1/2} = \text{constant}$. If absorption at the critical surface dominates $N = \text{constant}$.

The relation (3) yielding the burn mass as a function of energy absorbed may be compared with the condition for optimum gain (2). If these two relations are simultaneously satisfied the system will operate at peak gain even though the burn is incomplete, and for a fixed initial fibre diameter will give higher gains at lower energy. It is clear from the forms of equation (2) and (3) that only one such intersection can occur at (E^*, M^*) . In practice, of course, a second intersection may occur when the fibre is fully burnt through (Figure A4.1). The parameters (E^*, M^*) are important scaling parameters for the problem, and in a well designed scheme should approximate to the actual operating condition, and to the mass for maximum gain.

We may identify two practical situations in which the mass M^* is greater or less than the initial fibre mass per unit length, M_0 . If $M^* > M_0$ no intersection occurs, and if $M^* < M_0$, two optimised conditions can be achieved.

It is clear from equation (1) that the gain decreases with increasing burn mass, and that in consequence a system should operate with low mass. However, for small masses equation (1) is only valid provided the laser pulse times are not greater than the expansion time of the equivalent fibre. If this condition is violated heating occurs during the expansion (cooling) phase and the gain falls rapidly (Figure A4.2). This will occur at a cut-off mass M_c , which depends strongly on the temporal characteristics of

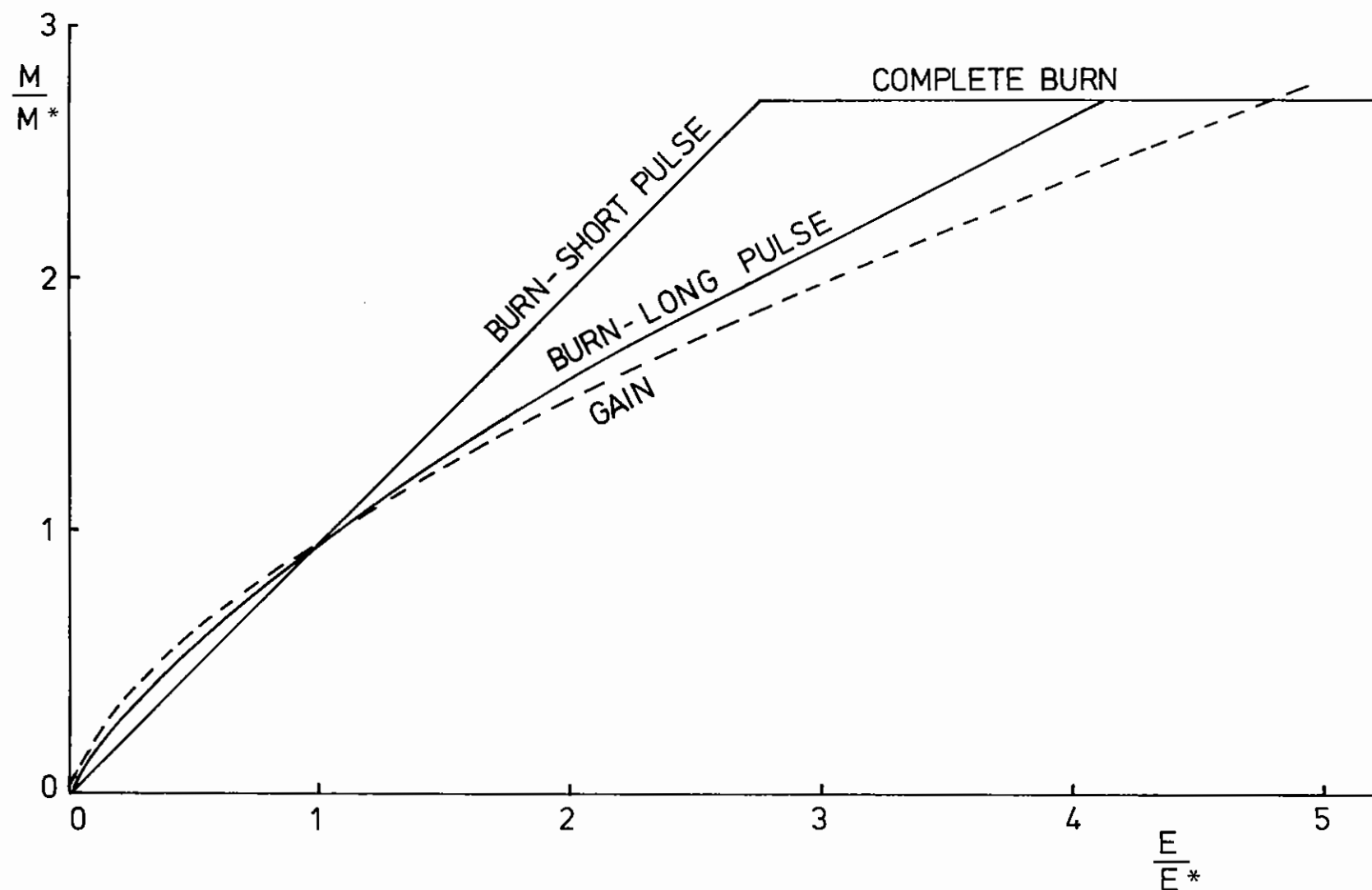


Fig. A4.1 Mass/energy plot for the model scaling laws for burn and gain. The mass and energy are plotted as ratios of the critical values (M^*, E^*). The initial fibre mass has been set at $2.75M^*$.

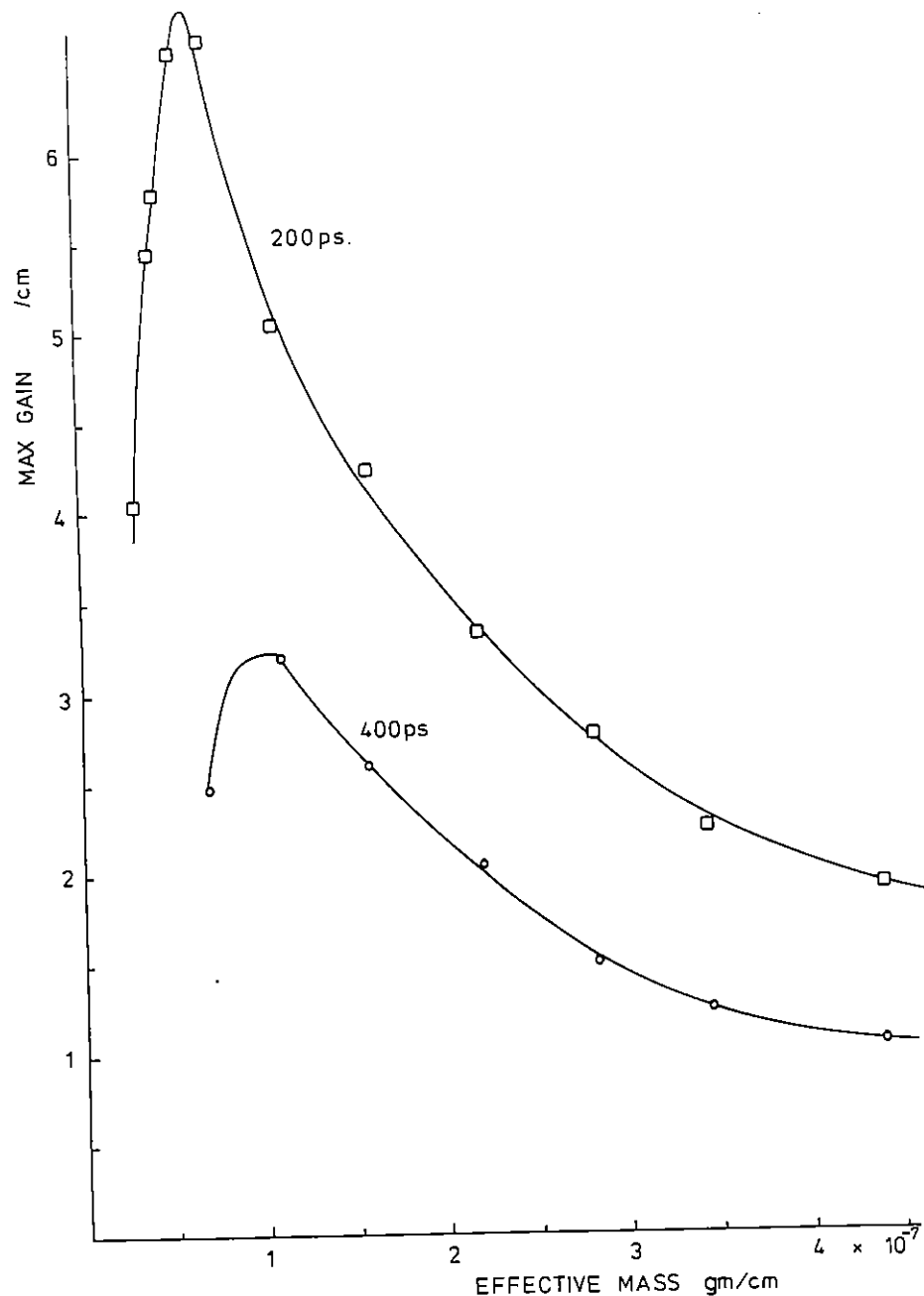


Fig. A4.2 Plots of peak gain versus effective mass for two laser configurations. The curve 200ps refers to one sided illumination by 1.06 μ m radiation with prepulse fraction 20% and delay 200ps. The curve 400ps describes four-sided illumination by 0.53 μ m radiation with prepulse fraction 20% and delay 400ps. The curve 200ps also gives a reasonable approximation to the case of four-sided illumination by 0.53 μ m radiation with prepulse fraction 20% and delay 200ps.

the laser pulse. For highest gain the pump laser characteristics should be set so that $M \approx M^* \approx N_c < M_0$, presenting a significant design challenge.

In practice it will be impossible to fully match the laser condition, within either focal spot or shot-to-shot variations, to the maximum gain condition everywhere and at all times. The gain decreases rapidly away from the neighbourhood of the peak gains as the energy is increased or decreased outside a window $E \pm 0.2 E$ where E is given by equation (2). On the other hand in the immediate neighbourhood of the peak the gain is relatively constant. Thus clearly the laser must operate within the gain window over its entire length, although not necessarily at its maximum gain value. This is achieved if the gain and burn curves closely match, ie in the vicinity of the intersection (E^*, M^*). From Figure A4.1 it can be seen that an energy window of $\pm 0.1 E$ allows a variation in the actual absorbed energy of between $\pm 0.2 E^*$ and $\pm 0.7 E^*$ for short or long pulses respectively: values which should be within the tolerance limits of working lasers. Under these conditions the gain along the fibre length will vary but will follow the variation as given by A4.2 for optimised gain as a function of burn mass. The optimal operating conditions (E^*, M^*) thus allows the system to naturally remain near peak gain and to take into account fluctuations in the laser power density. It must be remembered that these conclusions only hold if the variation in power, and therefore absorbed energy and effective mass are not too large, ie within the window specified above.

The model described in this section is clearly qualitative and no quantitative conclusions should be drawn. However, the general behaviour observed in computational calculations is in agreement with this picture. Unfortunately the optimal conditions (E^*, M^*) vary markedly with the laser pulse characteristics, and their values can only be deduced by a long and detailed set of computational runs. Figures A4.3 to A4.5 show three such sets of data for configurations of experimental interest. Figure A4.3 refers to the experiments carried out using 1.06 μ m radiation with one sided illuminations, pre-pulse ratio 20% and delay 200 ps. In this case $M^* \ll M_0$, and indeed reference to Figure A4.2 where the peak gain/mass curve for this configuration is plotted shows that $M^* < M_c$. This case, in which high gain was observed, is therefore expected to have a narrow gain window with a very

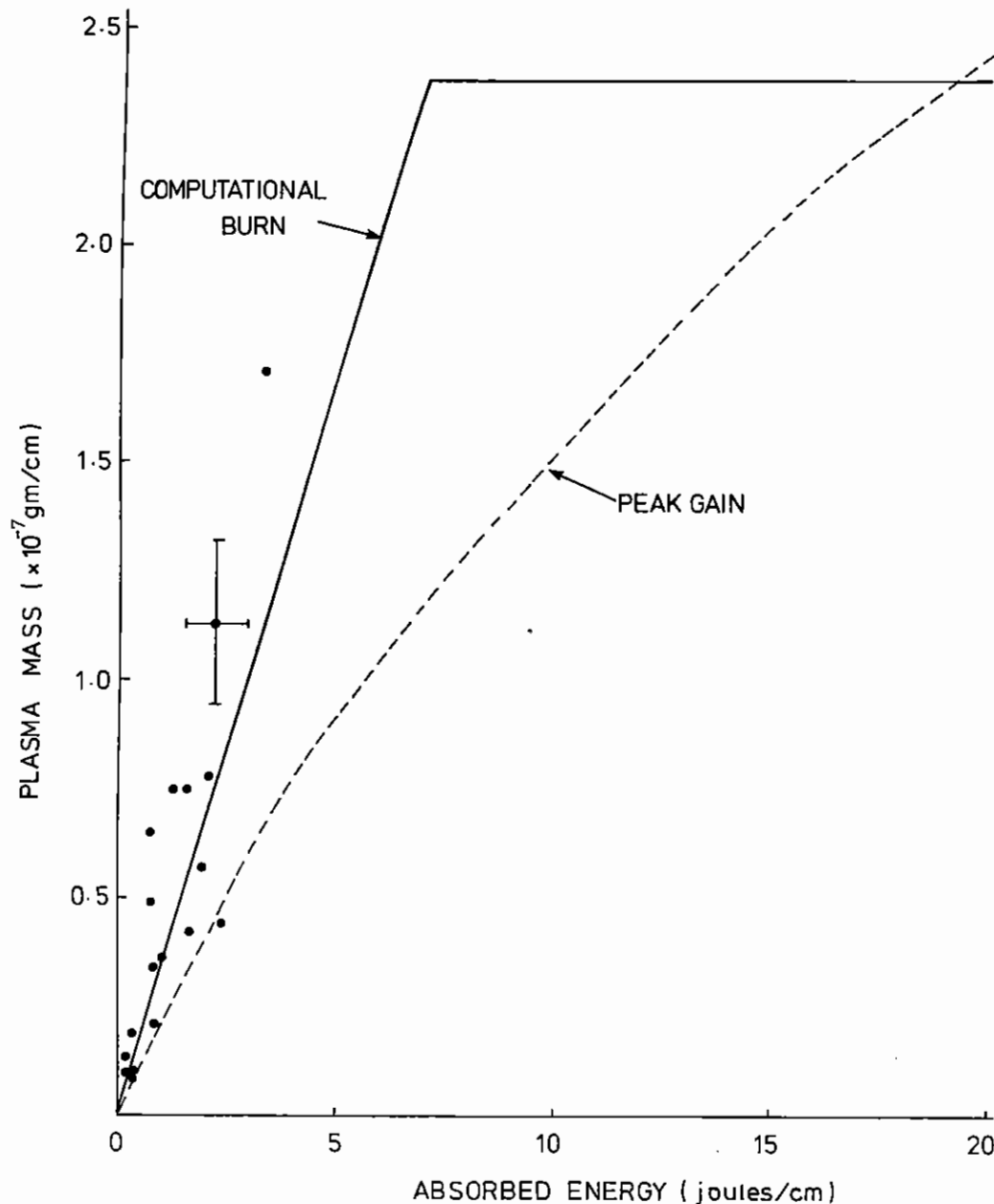


Fig. A4.3 Mass/energy plot for one-sided illumination of $4\mu\text{m}$ diameter fibres by $1.06\mu\text{m}$ radiation, prepulse delay 200ps and fraction 20%. Dots denote experimental measurements.

sharp cut-off at high energy as the burn and gain curves diverge. Such behaviour is consistent with the experimental data.

Figure A4.4 refers to the experiments carried out using $0.53\mu\text{m}$ radiation with four beam illumination, pre-pulse ratio 20% and delay 200 ps (section A4.2.3). In this case $M^* > M_0$ although optimal operation is possible at low mass, $M < M_0$. The experiments revealed no gain, and were consistent with non-optimised operation.

Figure A4.5 shows results corresponding to a system using $0.53\mu\text{m}$ radiation with four beam illumination, pre-pulse ratio 20% and delay 400 ps (A4.3). In this case the gain and burn curves lie close together and the system is well optimised. Unfortunately the gain cut-off occurs at rather high mass and correspondingly low gain. Thus although well optimised this system has too low gain for practical application. These conclusions are in good agreement with the experimental data obtained under these conditions (A4.3) which typically showed axial/transverse signals ratios about 1, and weak positive gain on some shots when account is taken of the time history of emission.

Examination of Figure A4.3 to A4.5 shows a marked difference between one and four-sided illumination burn characteristics. The one-sided case obeys the short pulse scaling law equation (3a) whereas both four-sided calculations follow the long pulse form (3b). This result, if consistent, represents a notable advantage for the latter case as it is clear from Figure A4.1 and Figure A4.5 that a considerably larger overlap of the gain/burn curves is possible in this case, and that the constraints along the fibre axis can be markedly reduced. Indeed in this case the restriction of tolerable power variation along the axis is likely to arise from the variation of timing of the gain onset and peak as the plasma mass varies rather than the actual value of the gain.

A4.2.2(a) Time Integrated Two Spectrograph Intensity Ratios

The standard diagnostic for gain in these experiments has been the single two spectrograph technique. Since such a method uses integrated signals

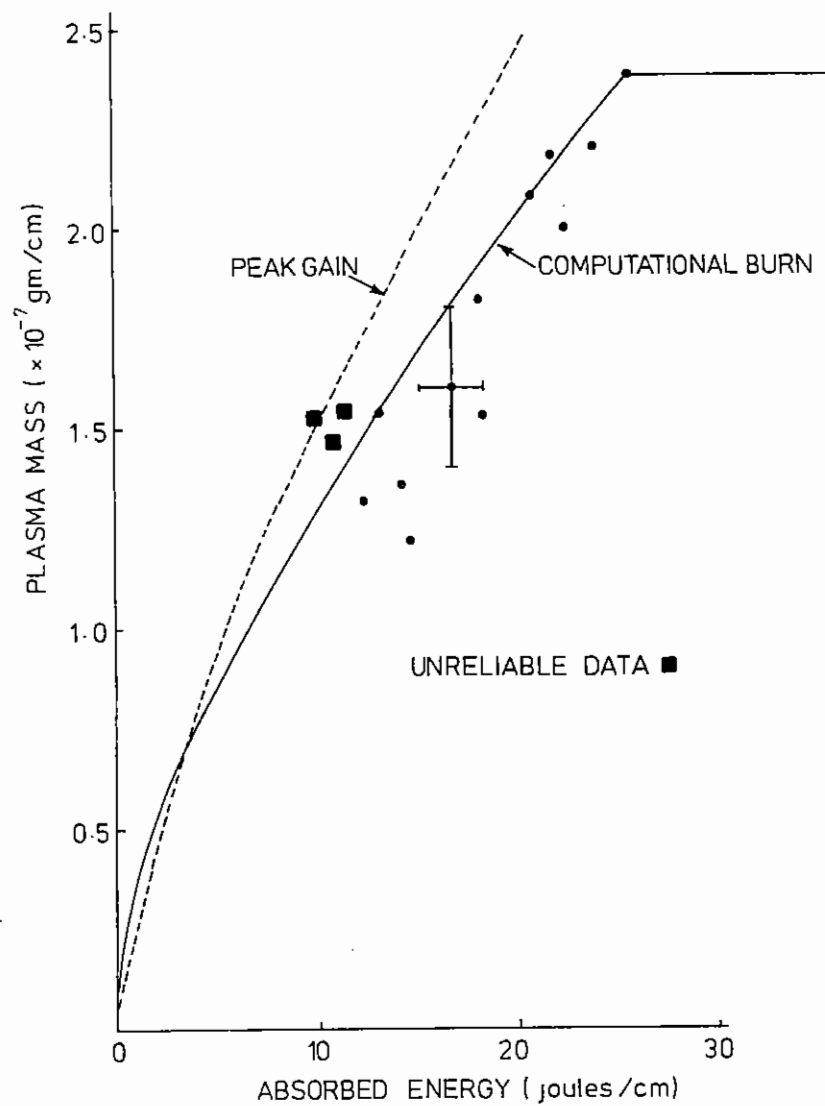


Fig. A4.4 Mass/energy plot for four-sided illumination of 4µm diameter fibres by 0.53µm radiation, prepulse delay 200ps and fraction 20%. Dots denote experimental points.

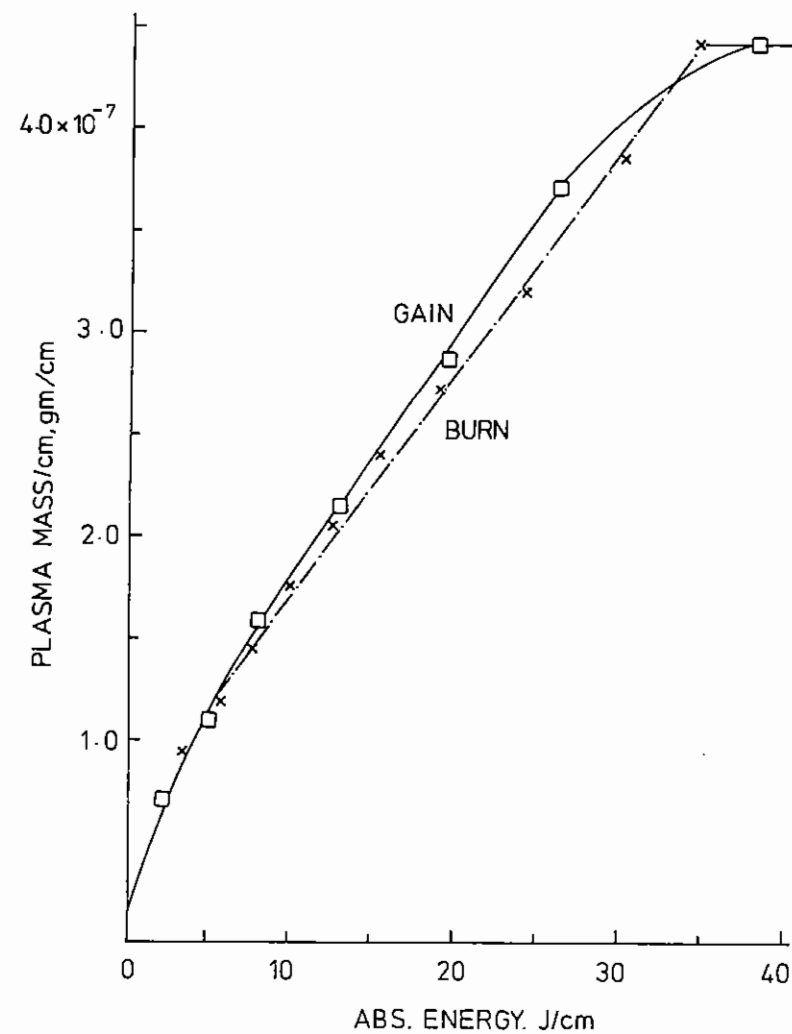


Fig. A4.5 Mass/energy plot for four-sided illumination of 5.5µm diameter fibres by 0.53µm radiation, prepulse delay 400ps and fraction 20%.

significant underestimates of peak gain will occur if the axial line is appreciably absorbed during emission. During the expansion the H_{α} line initially always has negative gain (ie absorption) before population inversion is established, due to the L_{α} line being optically thick. During this phase the axial H_{α} is effectively cut-off, the transverse H_{α} on the other hand is emitted without re-absorption. At the onset of gain the axial signal is thus markedly smaller than the transverse. This situation is of course reversed during the gain phase, however, as absorption occurs at high density but gain at lower, its effect is thus much stronger.

The reduction in observed signal amplitude due to this effect is readily estimated. Suppose a fraction α of the H_{α} signal is emitted before gain onset, and that the average gain amplification after gain onset $A = (\exp(Gl) - 1)/Gl$, then the ratio of axial to transverse signal is approximately

$$A' = A(1 - \alpha)$$

We note that an observation of gain requires $A > 1/(1 - \alpha)$.

During high mass expansions $\alpha \sim 30\%$, and $A > 1.4$ or $Gl > 0.8$, which for the short lengths of about 2 mm used in these experiments requires $G > 4 \text{ cm}^{-1}$, a value only possible at peak gain under these conditions (Figure A4.2).

At low mass $\alpha(\sim 10\%)$ is significantly reduced, and the gain ($Gl \sim 2$) is correspondingly high. In consequence the error in the measured gain is relatively small.

In order to estimate the consequences of this effect the code ITEST was modified to calculate the axial and transverse signals. Under the conditions of a typical run of the experiments reported in section A4.2.3 an axial/transverse ratio of 0.8 was found in reasonable agreement with the measured values (Figure A4.6). The computed peak gain length product was about 0.1.

Although the effective radius of the fibre in the earlier set of gain

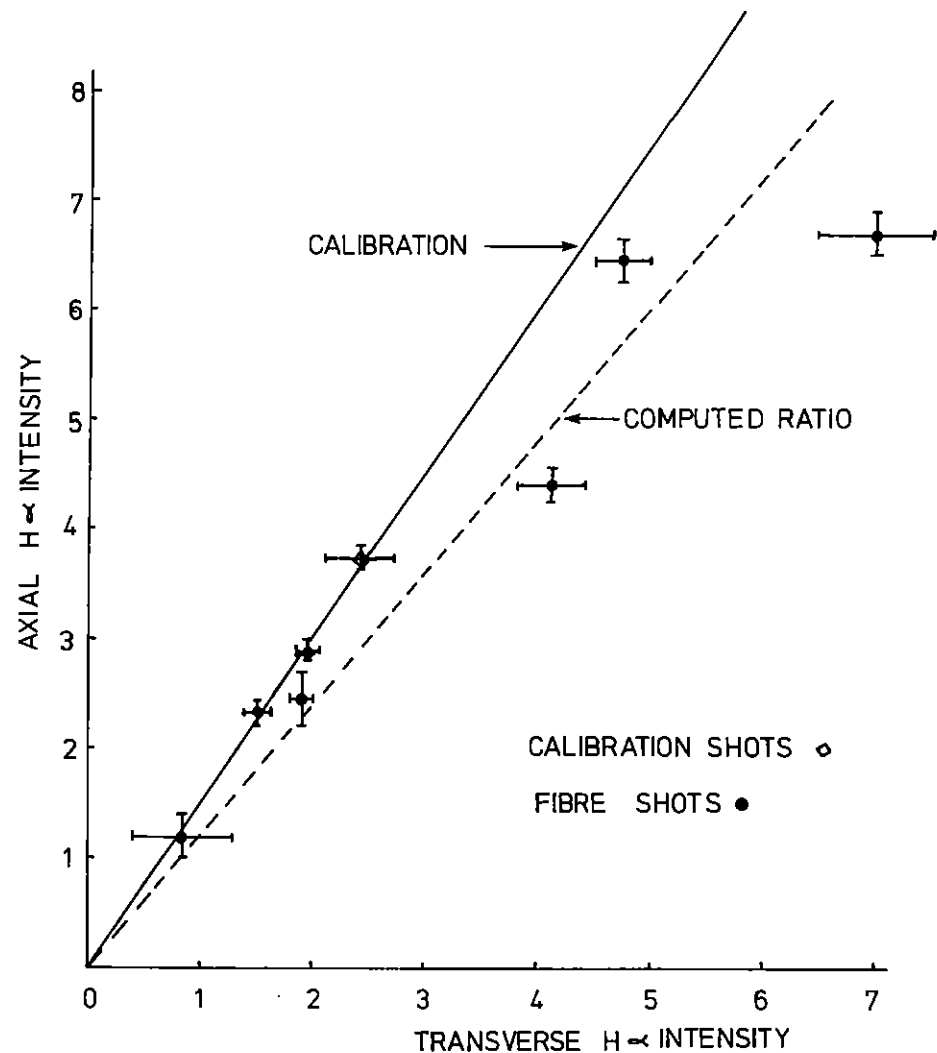


Fig. A4.6 Comparison of axial and transverse signal ratio measured experimentally and derived computationally for the experiments of Section A4.2.3

experiments (A4.3) was not measured an estimate from the computed data (Figure A4.5) was used to determine the input conditions for ITEST. In this case a signal ratio of about unity was found in agreement with the experiments. The computed peak gain length product in this case $G = 0.6$ corresponds to a gain of about 3 cm^{-1} .

A4.2.2(b) Mass/Energy Relationship Calculations

Preliminary calculations of the mass/energy relationship were reported last year (A4.3). These calculations were carried out using a version of the two-dimensional code POLLUX in which the bulk of the energy was dumped at the critical surface. It was noted that in these calculations for four beam geometry where reasonable cylindrical symmetry is to be expected there was marked disagreement between one and two dimensional calculations. We have now modified both one and two dimensional codes to include ray tracing so that an accurate representation of the energy deposition is obtained, and have resolved the discrepancy between the codes in favour of the one dimensional results.

Figure A4.7 shows the computed mass/energy relationship for the conditions of the experiments reported in section A4.2.3. Two dimensional calculations (+) were performed using a mesh 24×24 cells, and one dimensional () using 90 cells for an average ionisation $\bar{z} = 5$. It can be seen that there is excellent agreement between the codes provided the plasma mass is small. At large plasma mass the one dimensional code reproduces the experimental data but the two dimensional is in marked error. This error is traced to the small mesh size employed due to limitations of computer time. As the mass of the core becomes small (large plasma mass) the core occupies only one cell, and the representation of the ablation zone becomes similarly coarse. The ablation and heat transfer are therefore poorly represented with consequent error. In contrast the one-dimensional code being Lagrangian is not subject to this limitation. A test run with 80×80 mesh confirmed this conclusion.

The one sided illumination calculations (Figure A4.3) are also not as severely restricted by this problem as due to the absence of symmetry the

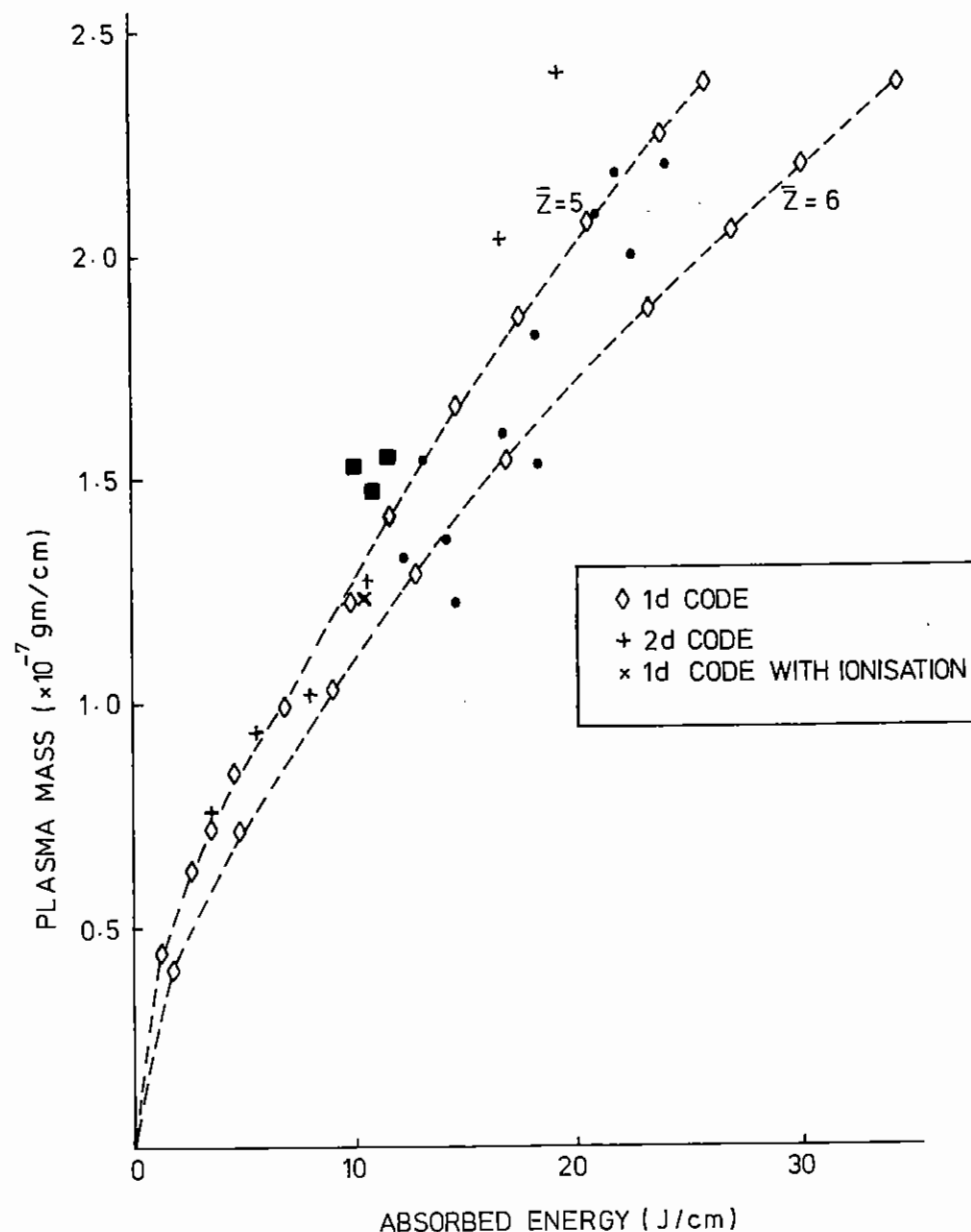


Fig. A4.7

Comparison of various numerical calculations of the mass/energy relation for the case given in Fig. A4.4. Good agreement between the 1d and 2d codes is obtained provided the plasma mass is not too large. The experimental data is consistent with an ionisation fraction $\bar{z}$ in the range 5-6.

core is not compressed during the heating phase, but expands and is displaced backwards retaining a significantly larger size than in the four sided case. We may again remark on the difference in scaling in these two cases, one-sided illumination following equation (3a) and four-sided equation (3b), in agreement with experiment.

The suitability of the one dimensional code to represent four sided heating considerably eases the computational requirement for identifying scaling laws and allows parameter variation. In figure A4.7 we show calculations for two values of z the average ionisation. $\bar{z} = 5$ has been preferred to account for the variation in the ablation zone for $\bar{z} = 0$ at the fibre to $\bar{z} = 6$ at the absorption surface, but both values are consistent with experimental data. A test run with a new version of LAGR including ionisation gave a value (x) of the burn mass intermediate between those of $\bar{z} = 5$ and $\bar{z} = 6$ (Figure A4.7).

A4.2.3 Experimental

The experimental arrangement was essentially the same as described in the last Annual Report (A4.5) with the target irradiated from four sides. The laser pulse was preceded by a prepulse ~ 300 ps before which had an intensity $\sim 20\%$ of the main pulse. The individual pulse durations were ~ 200 ps. As before the plasma was viewed along the axis and transversely by two grazing incidence spectrographs. The spectrographs were cross calibrated so that anomalous intensity ratios could be detected when using fibre targets. This intensity data for the H α line (182 Å) was to be compared with computational predictions under the same conditions.

Previous CVI experiments at both Hull and RAL have provided evidence for incomplete burn-through of carbon fibre targets (A4.6, A4.7). One of the main objectives of the last CVI experiment was to study the burn-through behaviour of fibre targets in much greater detail. Two independent techniques were employed to obtain this data. The first method was to use a charge collector in conjunction with a calorimeter array to determine the total number of ions produced. The second method was to use an

interferometric probe and to integrate the fringe shift data to determine the total number of electrons. It was also necessary to know the plasma column length accurately on each shot and a soft X-ray pinhole camera was employed for this purpose. The burn-through data would be compared with computational predictions. The ability to vary the time of the interferometric probe beam allowed "time resolved" measurements of density and radius which could also be compared with code predictions. All data presented here refers to fibres of initial diameter 4 μm .

A4.2.3(a) Calorimeter Array Data

The calorimeter array used was an improved version of the previously described device (A4.8) and had ten channels allowing a study of the energy distribution over an angular range of $\sim 24^\circ$. Each channel was constructed in a differential manner to subtract out the contribution of scattered laser light. The signals obtained were automatically logged and plotted using the GEC4080. A second ten channel array was used but behaved erratically, possibly due to thermal effects in the amplifier circuitry. When this array did work it gave signals consistent with those of the main array, suggesting that the expansion was reasonably symmetric (A4.5).

Figure A4.8 shows the signals obtained from the main calorimeter array, averaged over twenty shots on 4 μm fibres, and normalised to the peak. The line which has been drawn through the data is a best fit of the energy distribution predicted by a self similar model (A4.9).

$$\frac{E(\theta)}{E(90^\circ)} = \frac{k^5}{((k^2-1)\sin^2 \theta + 1)^{5/2}}$$

where $E(\theta)$ represents the energy per steradian at an angle θ to the fibre axis and k is the ratio of characteristic velocities parallel to the fibre axis and radially. Since the expansion is almost entirely radial we expect $k \ll 1$, the line drawn being for $k = 0.09$. This best fit value of k is to be compared with the code prediction $k \sim 0.105$. The assumption of a self similar expansion axially is unlikely to be valid and this probably accounts for the slight underestimate towards the free end of the fibre (angles less

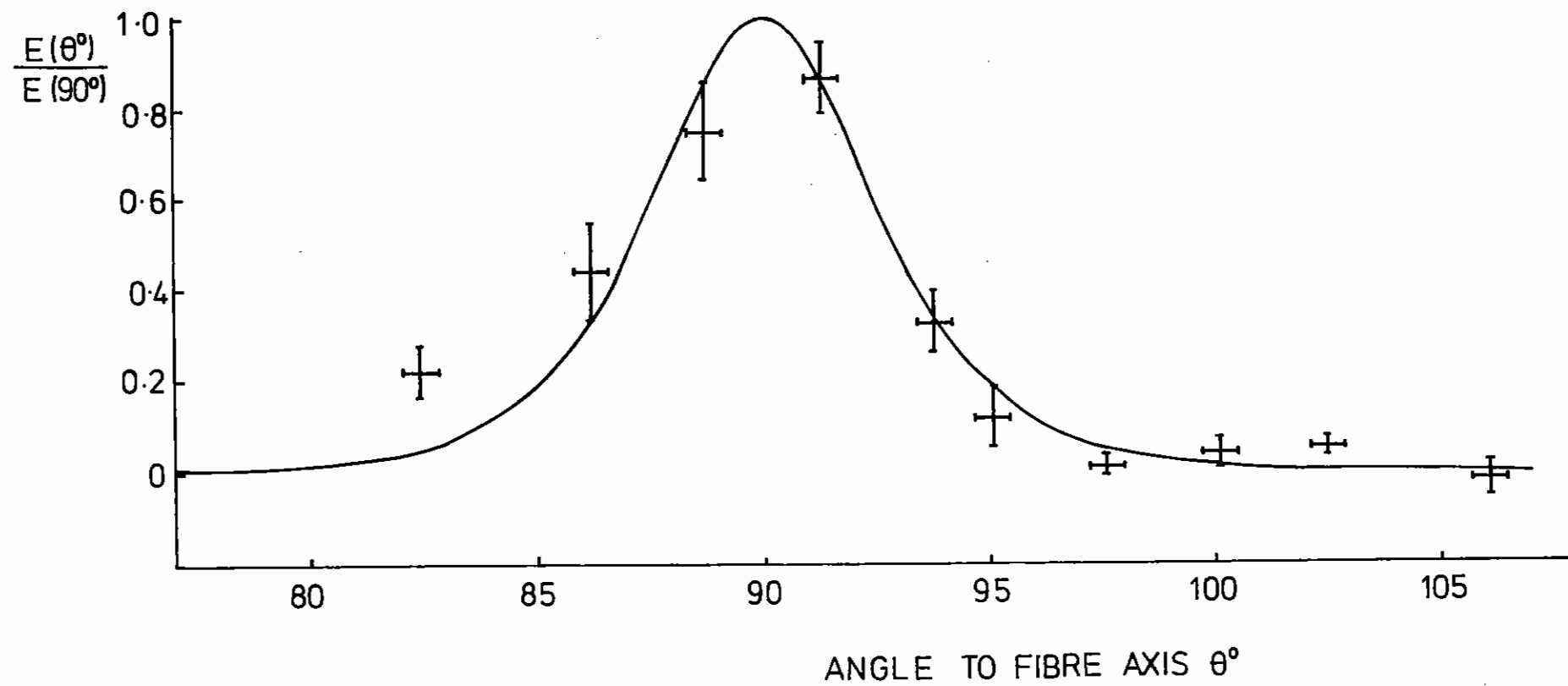


Fig. A4.8 Calorimeter array signals normalised to the peak and averaged over 20 shots on $4\mu\text{m}$ fibres. The curve is a similarity model prediction with $k = 0.09$ (see text).

than 90° correspond to the free end and angles greater than 90° to the fixed end). Having experimentally determined k and knowing $E(90^\circ)$ for any shot we can calculate the total energy coupled into the target (A4.5). The values we obtain agree to within $\sim 10\%$ with the much simpler peak $\times$ half width integration which is normally assumed.

A4.2.3(b) Charge Collector Data

A charge collector was placed in a position symmetrically equivalent to the calorimeter array position. The signals from this detector were displayed on a Tektronix 7844 oscilloscope and photographed. The signals were later digitized and integrated to give the total number of ions. By comparing the energy per steradian detected by the charge collector with the calorimeter signals the necessary correction for the "unknown" ion charge and expansion geometry can be found.

Figure A4.9 illustrates the variation of plasma mass with absorbed energy for $4\ \mu\text{m}$ fibres as measured experimentally. The solid line is the computational prediction of the plasma mass which is in good agreement with the measurements. We note that only one shot achieved complete burnthrough, requiring an absorbed energy of $\sim 26\ \text{J/cm}$. The dashed line in Figure A4.9 represents the computational prediction of the plasma mass and energy requirements needed to achieve peak gain. We note that the gain condition is not well matched to the burn-through.

By comparing the peak energy per steradian from the calorimeter array with the energy per steradian represented by the ion signal we can determine, from Figure A4.1, the value of θ for the charge collector (we assume the average ion charge $\bar{z} = 6$, which is consistent with the interferometry data). This value turns out to be only $\sim 3.7^\circ$ off peak. Hence we can use a similarity model to determine the total number of ions (A4.9). These values agree with the more directly determined numbers to within 10 - 15%.

Although the measured energies and ion numbers agree very well with a similarity model the comparison cannot be taken too far. If we examine the shape of the ion signals we find they are only qualitatively the same as the

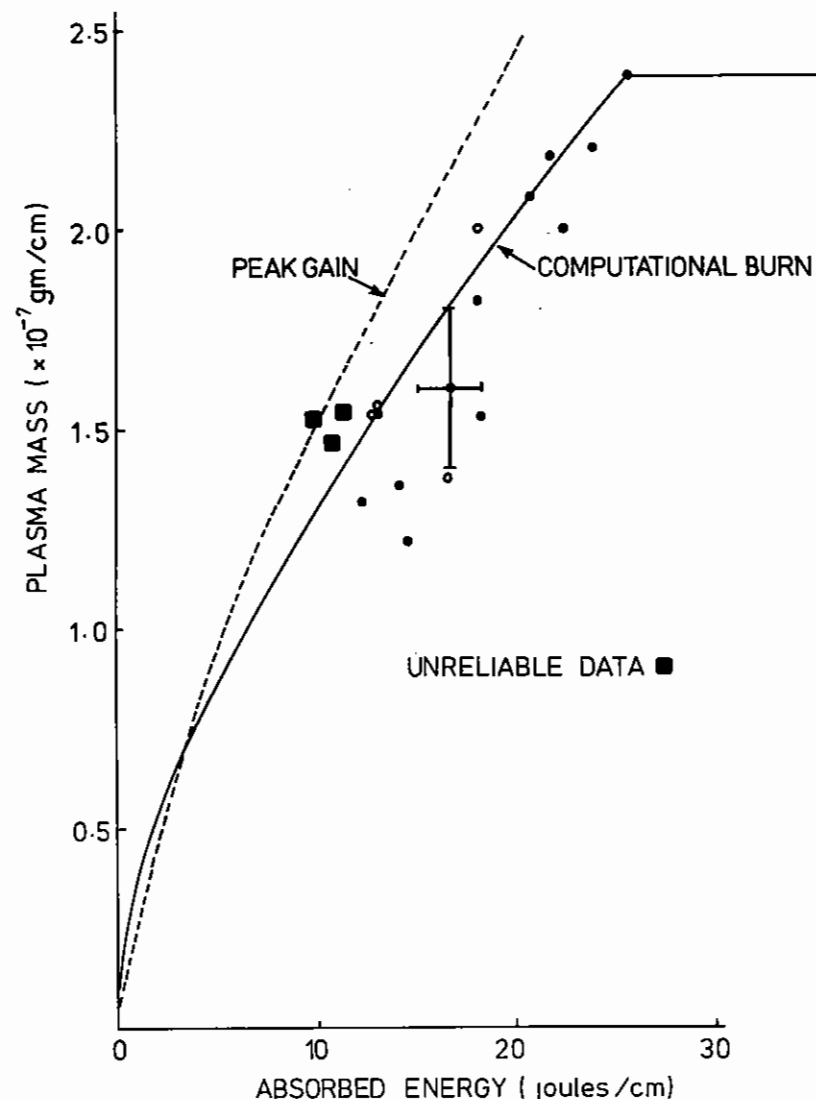


Fig. A4.9

Experimentally observed variation of plasma mass with absorbed energy ($4\ \mu\text{m}$ fibres). The solid points are charge collector data and the open points interferometry data. The solid line is the predicted burn-through behaviour and the dashed line is the predicted requirements for peak gain.

shape predicted by a similarity model. Figure A4.10 illustrates the ion signal obtained on a typical shot compared with the equivalent similarity profile. Both curves represent practically identical energies and numbers of ions but the time dependent behaviour is clearly different.

A4.2.3(c) Probe Beam Interferometry Data

A short pulse optical probe beam was used in conjunction with a Nomarski-type polarising interferometer to obtain interferograms of the plasma around the fibre. The expected plasma density at the time of interest, (1 ns after the main irradiation pulse) was quite low at 10^{19} cm^{-3} , so the probe wavelength had to be fairly long in order to obtain reasonably large fringe shifts. The wavelength chosen was 622 nm, which is obtained from the fundamental in a two-step conversion process.

The first step is to frequency-double the 1.05 μm , 200 ps probe pulse in a Type I CD*A crystal: this is in non-saturated mode, and produces approximately 30 mJ of green light in a pulse of ~ 150 ps. After filtering out the residual IR, this is down shifted by stimulated Raman scattering in an ethanol cell. The Raman process is about 80% efficient and results in a further pulse shortening by a factor between 3 x and 5 x, depending on intensity. The final pulse length is thus around 40 ps.

The experimental layout is as shown in Figure (A4.11). The probe light was concentrated onto the fibre with a 700 mm f1 lens, some difficulty being experienced in obtaining good enough uniformity over the necessary 2 mm field of view. The interferometer was set to give a fringe spacing in the object plane of $\sim 70 \mu\text{m}$, with the fringes tilted $\sim 25^\circ$ from the fibre axis. Alignment and focusing were done using He/Ne light, which is very close to the probing wavelength. The necessity of using the same optical path for fibre positioning meant that an extra 500 mm lens had to be introduced, once alignment was complete, to form an image of the required size at the camera: the magnification at the film plane was then 10 x.

The probe was timed to the peak of the main heating pulse using an S-1 streak camera, and interferograms were subsequently recorded at times of +

0.5 and + 1.0 ns on HP5 film. One of these is reproduced in Figure A4.12, and exhibits fringe shifts of up to two fringes in the region adjacent to the fibre.

One problem encountered was a partial loss of coherence in the beam across the object plane: this led to regions of fringe confusion and made analysis of the interferograms difficult. The cause was traced to beam breakup in the ethanol cell, resulting from too high a green intensity, and was largely cured by filtering out some of the second harmonic before focusing it into the cell.

The interferometric probe beam was used on four of the shots on fibres. The fringe shift data obtained at a particular position along the length of the fibre can be integrated directly to give the number of electrons per unit length. Knowing the length of the pinhole camera image and assuming uniformity along the length of the plasma column we may then determine the total number of electrons. In all four cases this is within 15% of the fully stripped condition $\bar{z} = 6$. Assuming $\bar{z} = 6$ we may plot the interferometry data on Figure A4.9 (shown as open points). This data is clearly in good agreement with the charge collector data.

The fringe shift data can be Abel inverted to obtain electron density profiles. An example of two density profiles is shown in Figures A4.13. The two profiles are shown together because they were obtained under almost identical conditions but at different times (0.5 ns and 1.0 ns with respect to the main pulse peak). Also shown are computational predictions for the two profiles which we see are in remarkably good agreement. In particular the small bump in electron density is quite well described.

Although the other two interferograms were obtained under slightly different conditions we can, nevertheless, use the extra data to gain some idea of the time dependent behaviour of the plasma. Since we expect to generate an inversion at a radius of $\sim 100 \mu\text{m}$ it is of interest to determine the density at this radius and compare the values with computer predictions. It is necessary to assume some average conditions for the energy input to the code. Figure A4.14 shows this data together with the computer predictions.

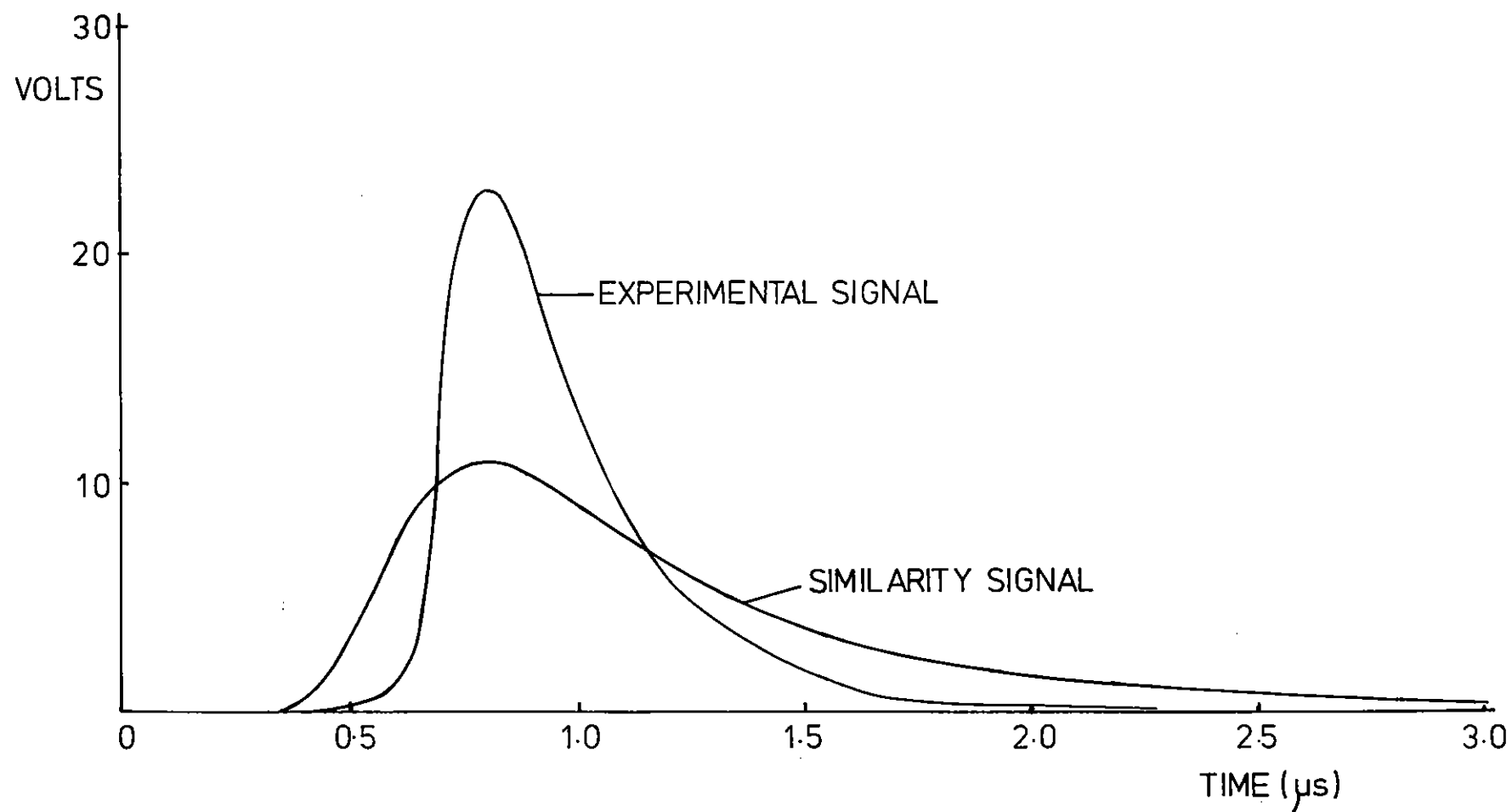


Fig. A4.10 Comparison of the experimentally observed ion signal and that predicted by a similarity model for shot 06 24/11/83.

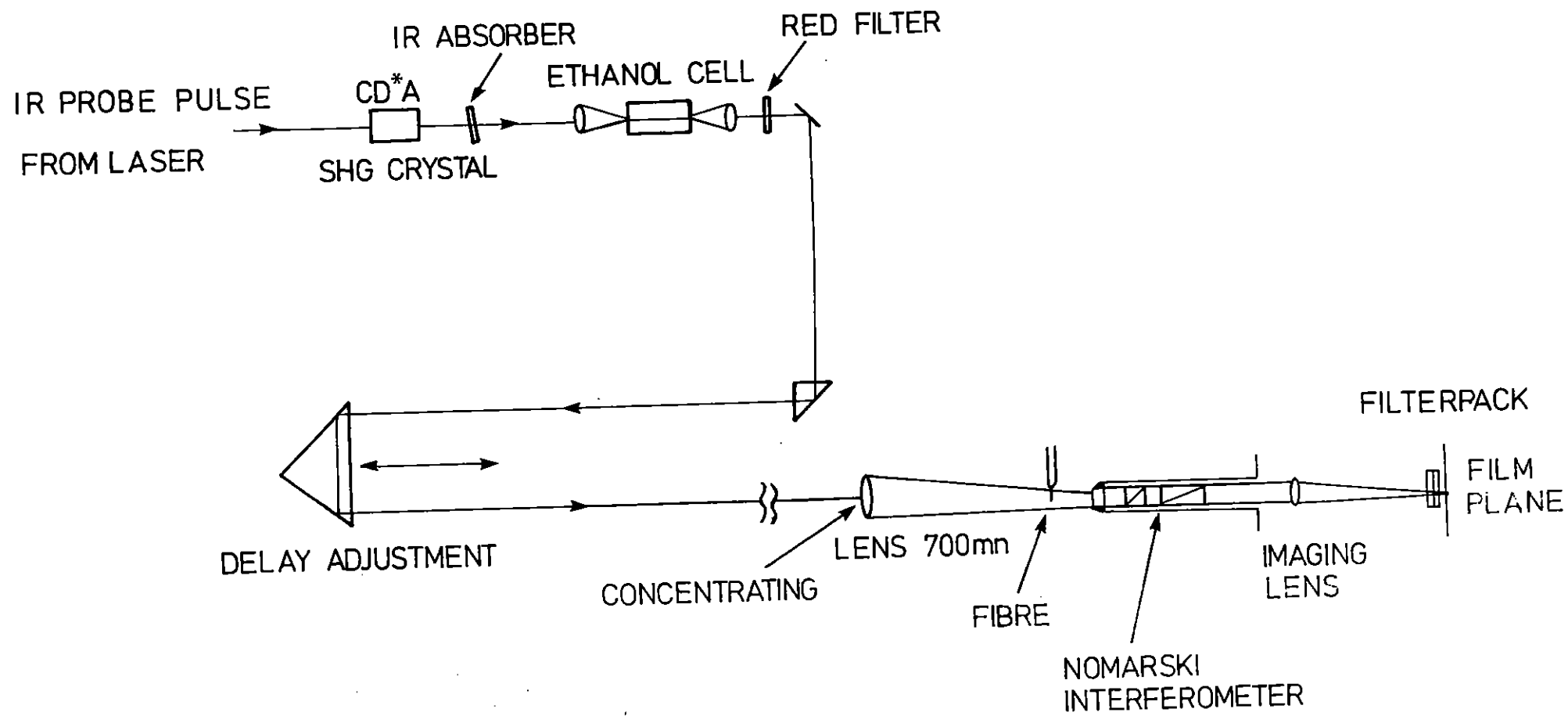


Fig. A4.11 Probe beam layout for carbon fibre interferometry.

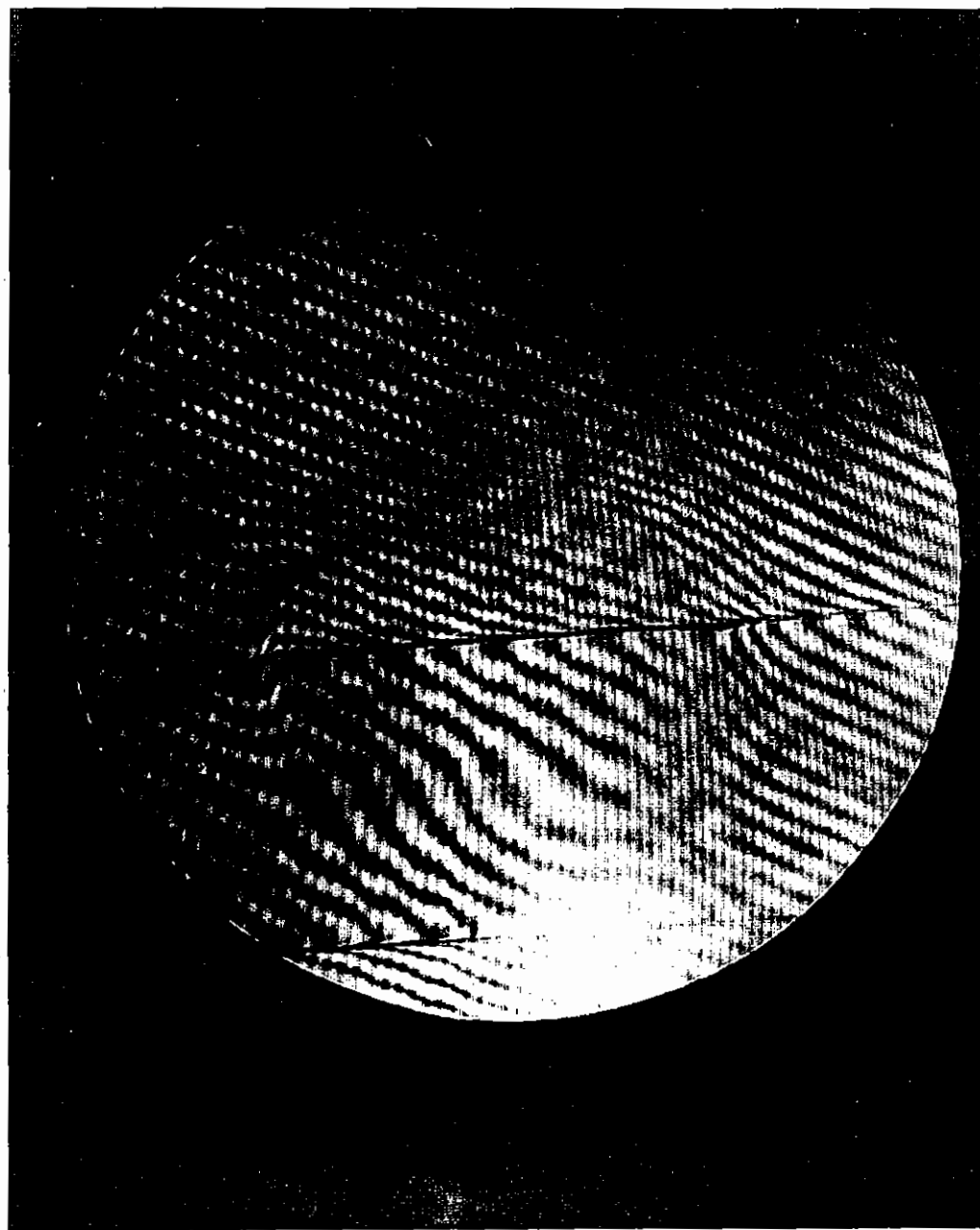


Fig. A4.12 Interferogram obtained on shot 20 24/11/83 at a time of 1.0ns after the peak of the main pulse.

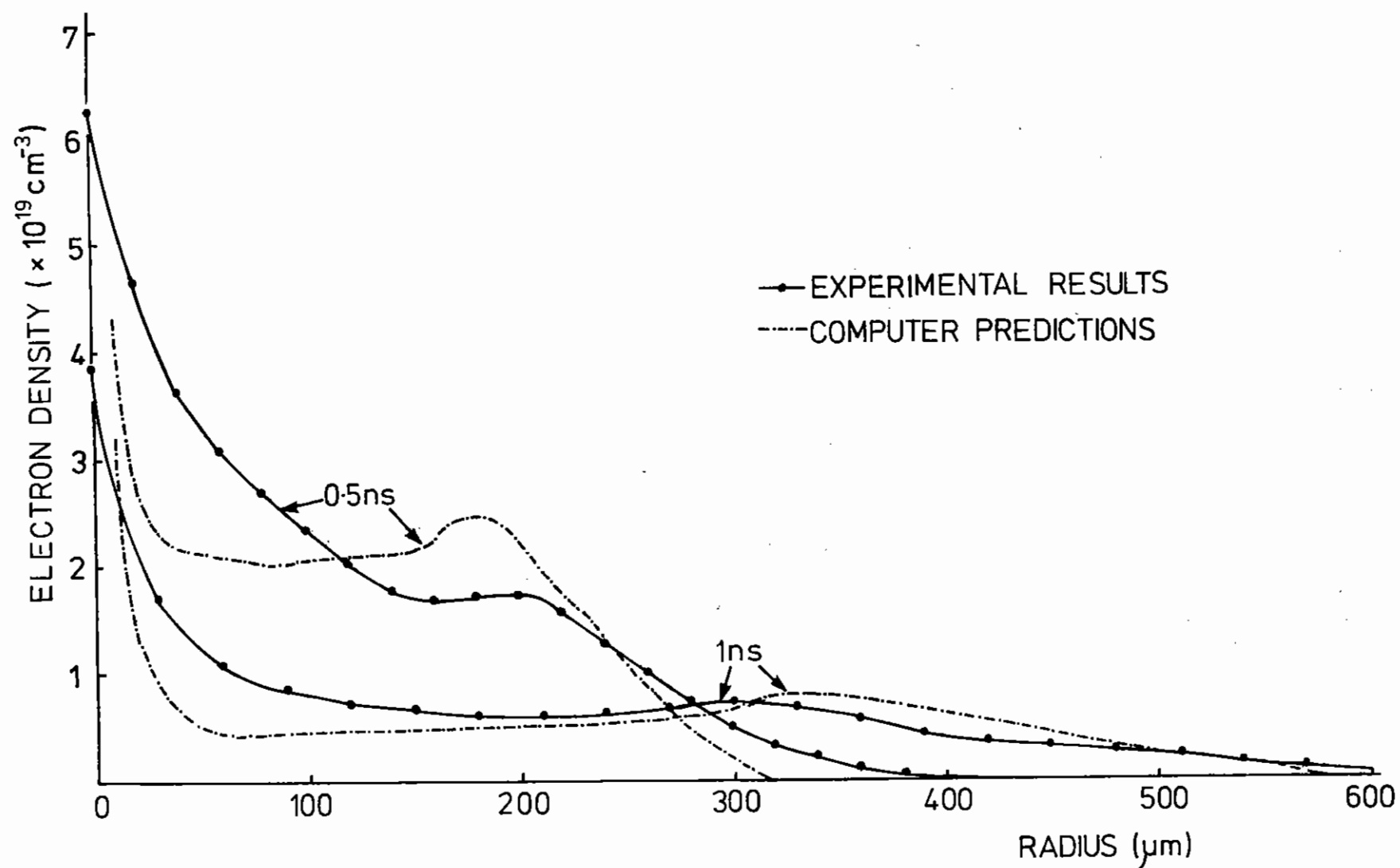


Fig. A4.13 Electron density profiles for shots 14 and 20 on 24/11/83 at times of 0.5ns and 1.0ns after the main pulse peak. The absorbed energy was $\sim 13\text{J/cm}$ in each case. For comparison the computer predictions for these times and conditions are also shown.

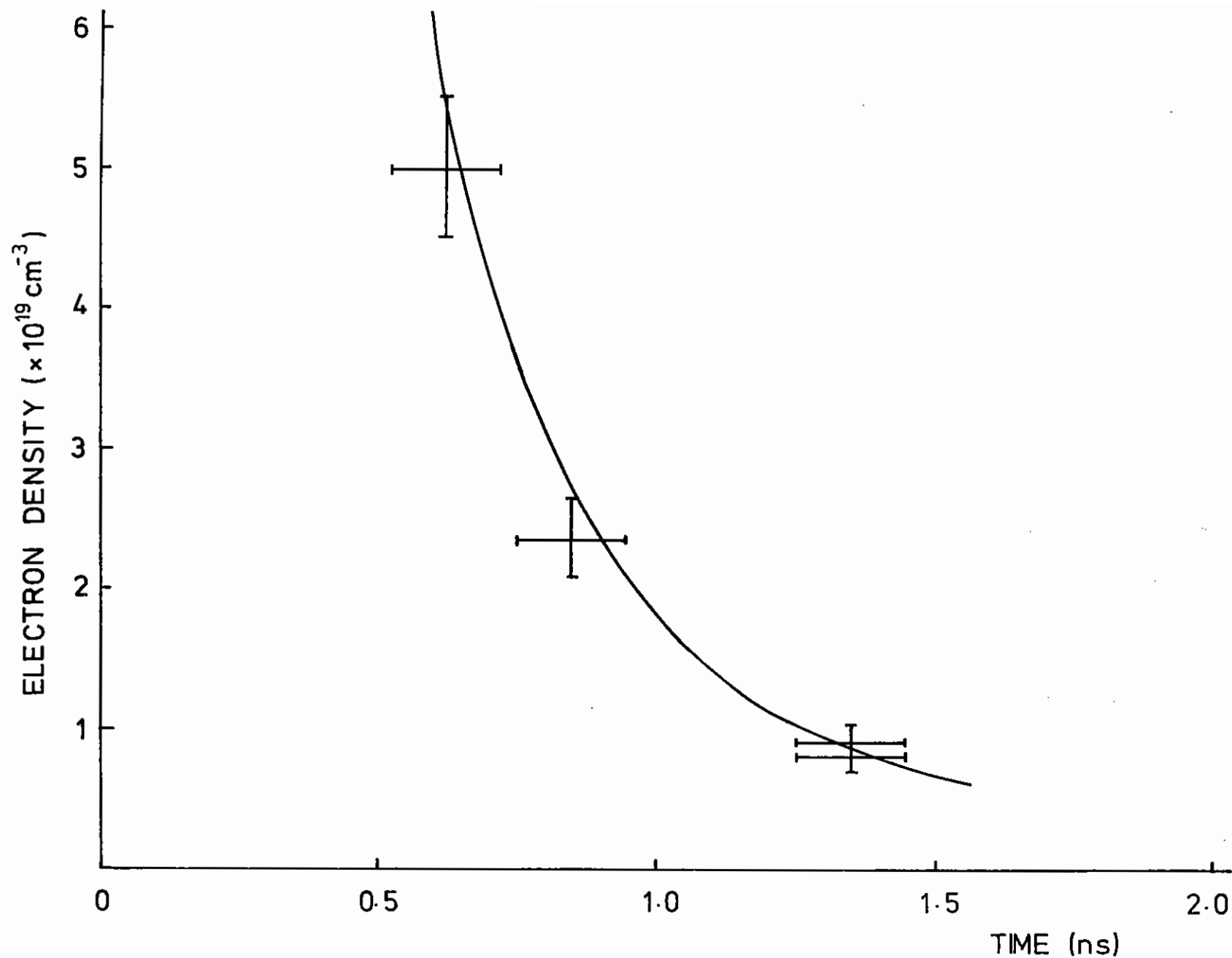


Fig. A4.14 Electron density at a radius of $100\mu\text{m}$ as a function of time. The curve is the computer prediction for laser conditions typical of the experimental data. Times are with respect to the prepulse onset.

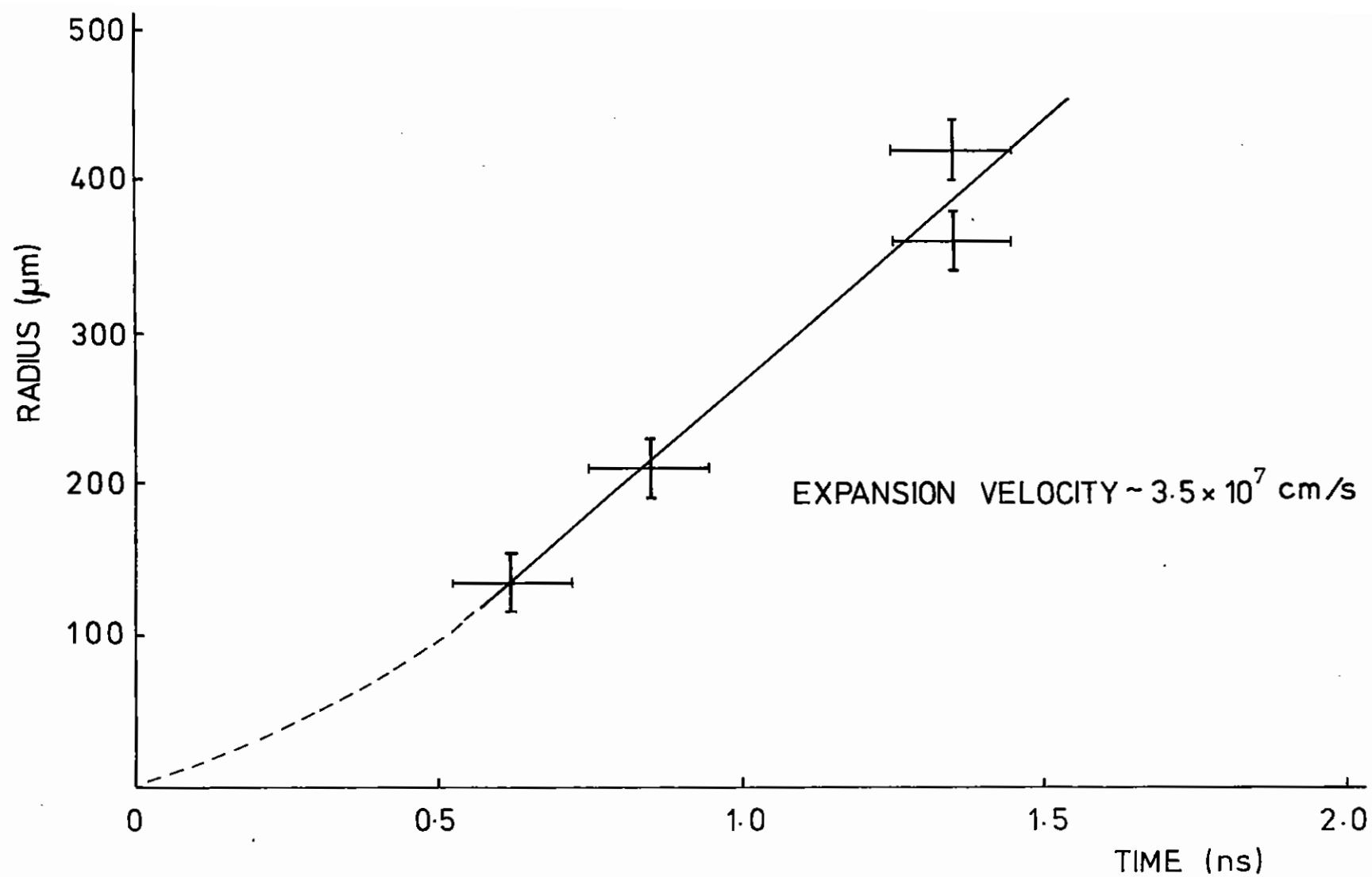


Fig. A4.15 Plasma radius as a function of time. The line is a visual best fit. Times are with respect to the prepulse onset.

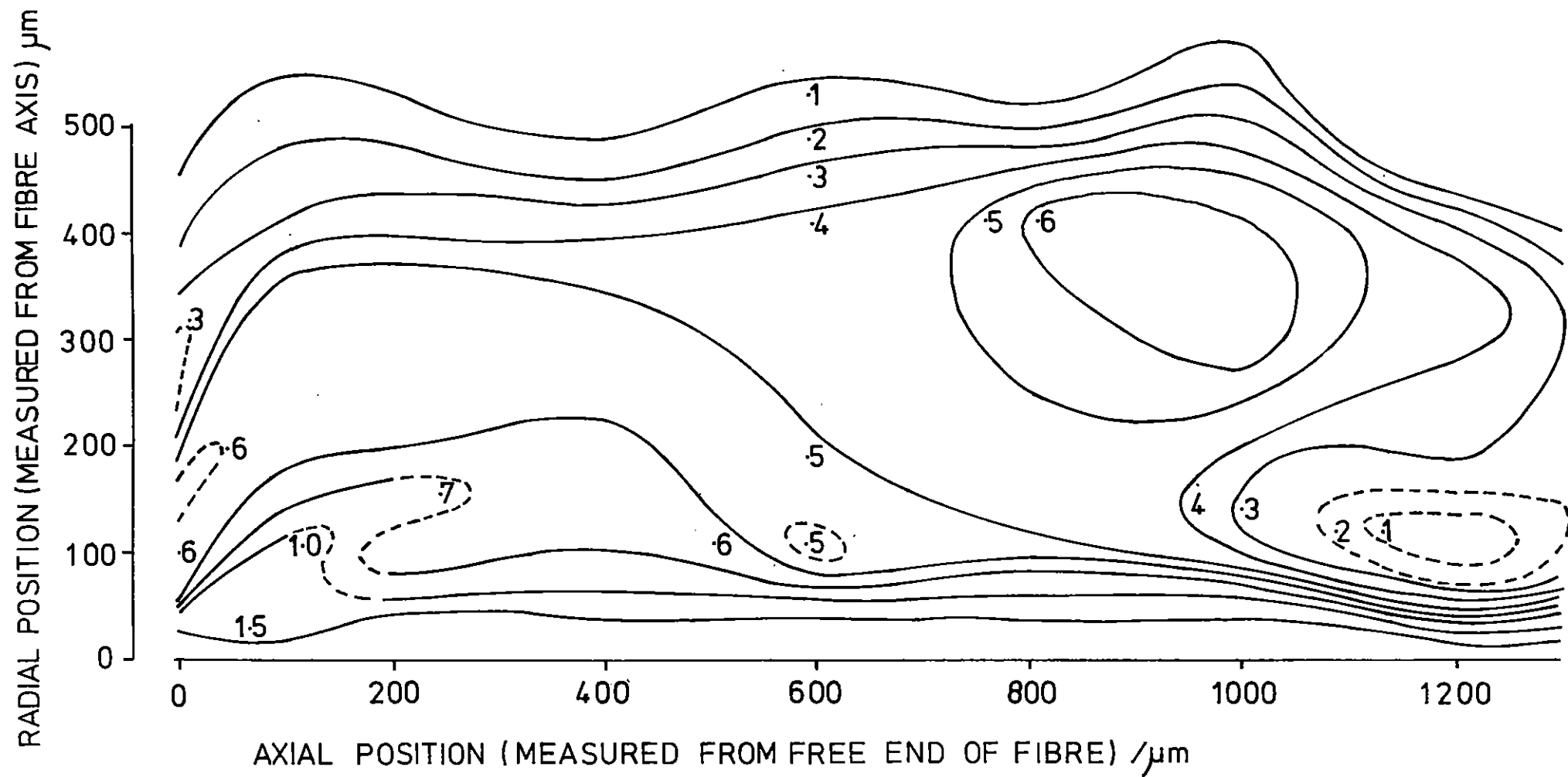


Fig. A4.16 Electron density contours for shot 20 24/11/83. The contours are in units of 10^{19} cm^{-3} and the data is at a time of 1.0ns after the main pulse peak (see Fig. A4.11).

The errors in the time are due to uncertainty in the prepulse onset time.

In Figure A4.15 we show the plasma radius as a function of time. We have defined the radius as that distance which encloses 63% of the total number of electrons within the plasma (corresponding to the $1/e$ radius of a Gaussian profile). An expansion velocity $\sim 3.5 \times 10^7 \text{ cm s}^{-1}$ is indicated which agrees with the measured ion velocities of $\sim 4 - 5 \times 10^7 \text{ cm s}^{-1}$ (velocity represented by peak of ion signal). Since there is no unique definition of radius or of ion velocity the agreement can be regarded as good.

It is possible to match computer predictions of the radius to the data in Figure A4.15 by a suitable choice of definition of radius. The problem here is that there is no clear distinction between cold core and plasma so the number of electrons out to a given distance is somewhat uncertain. The line drawn in Figure A4.7 is simply a visual best fit and does not represent a rather arbitrary prediction.

The major assumption is comparing the interferometry data with the charge collector data was that the plasma was uniform along its length. To test this assumption one of the interferograms has been scanned at several positions along the plasma column and a set of density contours built up. This data is shown in Figure A4.16. Apart from a relatively small density bump at $\sim 900 \text{ }\mu\text{m}$ and a large depression at $\sim 1200 \text{ }\mu\text{m}$ the contours are reasonably uniform. In fact these two features may not be real anyway since they correspond precisely to regimes of poor fringe visibility on the interferogram (Figure A4.12). In any case the uniformity is good over a length of at least 0.5 mm and the assumption of uniformity in the data analysis seems justified. It is a little less clear whether the observed non-uniformities have any significant effect on the observation of gain.

A4.2.3(d) Spectroscopic Data

A number of problems marred the spectroscopic data obtained during the experiment and it was necessary to completely disregard the initial calibration. As a result only a limited amount of data was usable. This

was made worse by the fact that single shot spectra from fibre targets proved impossible and this effectively halved the shot rate. Figure A4.17 presents all the useful spectroscopic data for the H_{α} line ($182 \text{ }\text{\AA}$). It is clear that all shots on fibres fall either on the calibration line or significantly below it, indicating zero gain or absorption. Referring to Figure A4.9 we note that the conditions used were not well matched to the gain requirements anyway but, nevertheless, some gain does occur. However, during the history of the H_{α} emission, strong absorption also occurs and a time integrated measurement will be dominated by this absorption, unless the gain is high. The dashed line in Figure A4.17 illustrates the expected intensity ratio, integrated over time, for the particular conditions used. This prediction is at least qualitatively in agreement with the ratios actually observed.

A4.2.4 Discussion

Although the experiment was rather disappointing in that gain was not observed it did provide a considerable amount of data in excellent agreement with code predictions. This suggests that if the correct conditions can be met then gain will be observed. Indeed it would appear that the experiments described in the last Annual Report (A4.5) were probably better matched and so some gain was seen. Note that in Figure A4.9 the gain condition lies to the left of the burn-through curve whilst the equivalent figure in the last Annual Report shows it to the right, so that it ought to be possible to achieve good matching using some 'mid-ground' conditions.

To improve the detectability of gain a number of routes could be chosen. Increasing the length of the plasma column is the simplest approach and the design of good quality line focus optics to give a length of $\sim 1 \text{ cm}$ should be investigated. The gain region within the plasma is fairly localised at a radius of $\sim 100 \text{ }\mu\text{m}$ so that space resolved measurements using a continuum source to probe the plasma may be a useful approach. Ultimately one could consider time resolution with all the attendant experimental difficulties that it entails.

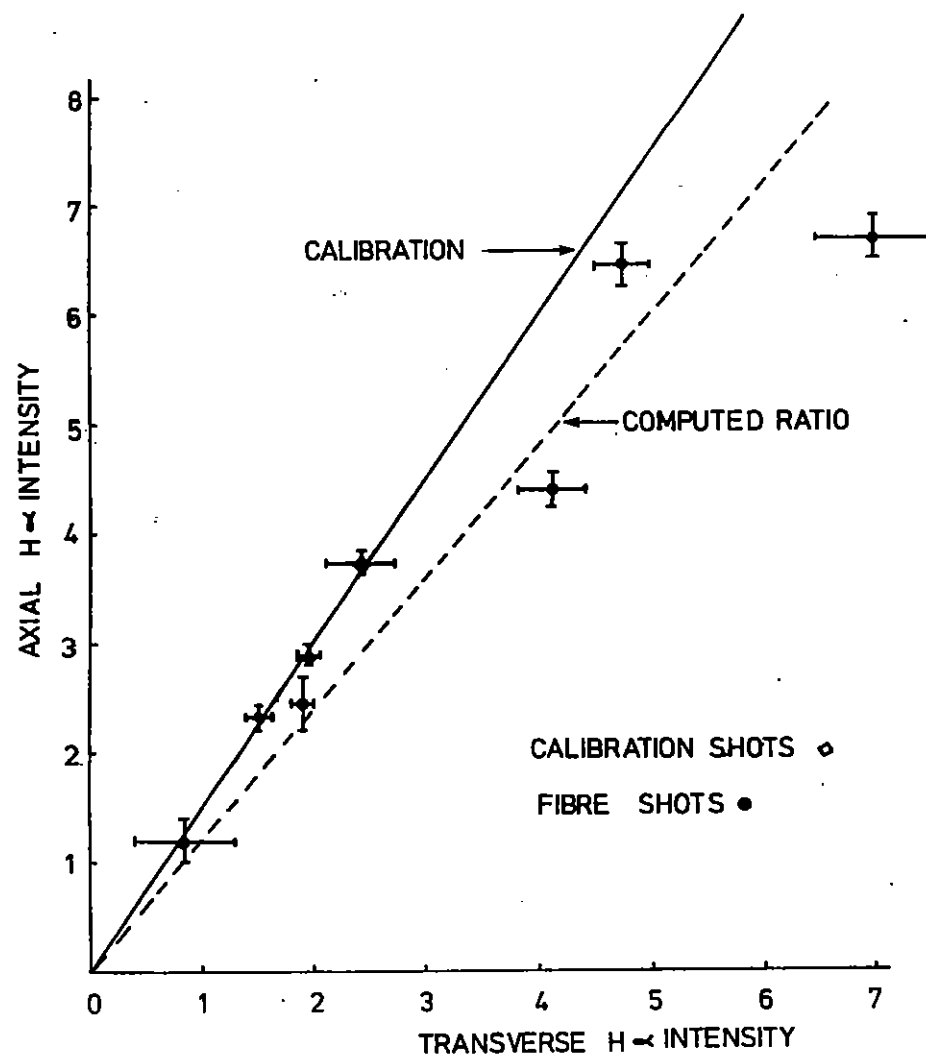


Fig. A4.17 $H\alpha$ intensities (in arbitrary units) observed on calibration shots and fibre shots. The computed ratio indicates that time integrated measurements should show absorption, in agreement with the experimental data.

A4.3 Soft X-ray Population Inversion Using Resonant Photo-Excitation

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A4.3.1 Introduction

Some preliminary experiments related to photo-pumped X-ray laser schemes have been carried out. These consisted of spectroscopic measurements from doubly irradiated thin foil targets in which pumping and pumped plasmas were produced in close proximity on opposite sides of the target.

The analysis of the results has concentrated on measurements related to the Sodium-neon scheme (A4.10) and this is described in the following section. The scheme involves the pumping of the $1s^2 \ ^1S - 1s4p \ ^1P$ transition in Ne IX by the $1s^2 \ ^1S - 1s2p \ ^1P$ resonance line in Na X. Suitable targets were made using the technique of ion-implantation. Neon was implanted on one side of a thin aluminium foil and sodium fluoride was coated on the other side. A 100 ps laser pulse at $2 \times 10^{14} \text{ W cm}^{-2}$ was used to irradiate the neon side, and 330 ps later the sodium fluoride side was irradiated with a similar pulse. X-ray spectroscopy was used to measure the state of ionisation of the neon, and the inferred electron temperature was compared with a one-dimensional computer simulation. The spectral brightness of the pump radiation was also estimated.

Measurements were made of the spectral brightness of chromium and tellurium which may prove useful pump sources in the 1 keV region. Tellurium emits strong continuum radiation in this region and so avoids the problem of wavelength matching which is critical with narrow line pump sources.

A novel X-ray laser scheme has been proposed by Lunney (A4.11) involving the resonant photoexcitation of autoionising states of lithium-like ions. For example, the Lyman- α line in F IX may be used to pump the $1s^2 2s \ ^2S_{1/2} - 1s2s \ (^3S) \ 3p \ ^2P$ transition in F VII leading to gain at 105.6 Å. Preliminary experiments relating to this scheme consisted of measurements of the pump intensities available in the F IX Lyman- α line and in the alternative pump

sources mentioned above.

A4.3.2 The Sodium-Neon Scheme

The plasma conditions required for obtaining maximum gain in resonantly photo-pumped neon IX have been considered theoretically by Apruzese et al (A4.10). The Na X $1s^2 1S - 1s2p 1P$ line at 11.0027 Å can be used to pump the Ne IX $1s^2 1S - 1s4p 1P$ line at 11.0003 Å and so give population inversion on the 2-4, 3-4 and 2-3 singlet transitions in Ne IX at 58, 230 and 82 Å respectively. It has been predicted that a gain of $\sim 100 \text{ cm}^{-1}$ can be obtained for the Ne IX $1s2p 1P - 1s3d 1D$ line if the neon plasma has an electron temperature of 65 eV, an electron density of 10^{21} cm^{-3} and is photo-pumped with radiation flux which is equivalent to a 227 eV blackbody. The required plasma conditions can be obtained in the expansion phase, following irradiation, of a laser produced plasma. If the neon plasma can be placed close to a laser produced plasma containing sodium then the required pumping flux can also be obtained. It is necessary to prevent over-population of the lower lasing level by photon trapping of the $1s^2 1S - 1s2p 1P$ line. Hagelstein (A4.12) has calculated that the opacity of this line should be less than 5, which implies that

$$n(\text{Ne IX}) \times r \leq 7.5 \times 10^{15} \text{ cm}^{-2}$$

where $n(\text{Ne IX})$ is the density of helium-like neon and r is the shortest dimension of the neon plasma.

The pumping and pumped plasmas should be placed in close proximity so that a large solid angle of pumping radiation is coupled to the pumped plasma. Various schemes have been suggested to produce the two plasmas. Hagelstein (A4.12) has suggested enclosing neon gas in a thin cell, the sides of which can be irradiated to form the pumping plasma. It has also been suggested that the neon plasma can be prepared by the laser irradiation of a thin layer of neon frozen on a solid substrate (A4.13). In the method adopted here neon was ion-implanted on one side of a 3 µm aluminium foil which is transparent to the pumping radiation at 11.00 Å. Sodium fluoride was coated on the other side of the foil as shown in Figure A4.18. The neon was

implanted at 80 keV which gives a range of 0.14 µm and a straggling distance of 0.046 µm (A4.14). This range was chosen since for a 100 ps, 0.53 µm laser pulse at $10^{14} \text{ W cm}^{-2}$ the ablation depth is $\sim 0.1 \text{ µm}$ (A4.15). Assuming that $\sim 50\%$ of the neon is He-like then the above opacity constraint dictates that the surface dose of neon should be less than $1.5 \times 10^{16} \text{ cm}^{-2}$. Some targets were prepared at this doping level and also at $6 \times 10^{16} \text{ cm}^{-2}$ to give more intense X-ray line emission; only the targets at the higher doping level were used here. A 0.1 µm layer of sodium fluoride was used so that the ablation front would penetrate through the sodium fluoride and yet ensure that the H-like and He-like resonance lines are optically thick. The targets were mounted in a brass ring to give laser and diagnostic access to both sides.

The expansion and cooling of the plasma on the neon side of the target was modelled using MEDUSA. Figure A4.19(a) shows the MEDUSA predictions for spatial variation of the plasma conditions at the end of a 100 ps laser pulse when a plane aluminium target is irradiated at $10^{14} \text{ W cm}^{-2}$. The region of the plasma corresponding to the neon-implanted zone is indicated and has reached an electron temperature sufficient to give H-like and He-like neon. Figure A4.19(b) shows the predictions for the plasma conditions 330 ps after the peak of the laser pulse. At the time the plasma containing the neon has moved 50 µm away from the target surface, the electron temperature is $\sim 100 \text{ eV}$ and the electron density is $\sim 5 \times 10^{20} \text{ cm}^{-3}$. These plasma conditions are close to the predictions of Apruzese et al (A4.10) for maximum gain on the 82 Å laser line; so for the targets considered here this is the appropriate time to apply the X-ray pumping pulse.

The targets were irradiated using the frequency doubled output of two beams of the VULCAN laser system. On the neon side $\sim 6 \text{ J}$ of laser energy in a 100 ps pulse was focused to a 220 µm diameter spot giving a laser irradiance of $2 \times 10^{14} \text{ W cm}^{-2}$. The focal spot size was measured by X-ray pinhole photography. A 10 J pulse of the same duration was focused to the same spot size directly behind the laser spot on the neon side and was timed to arrive 330 ps later.

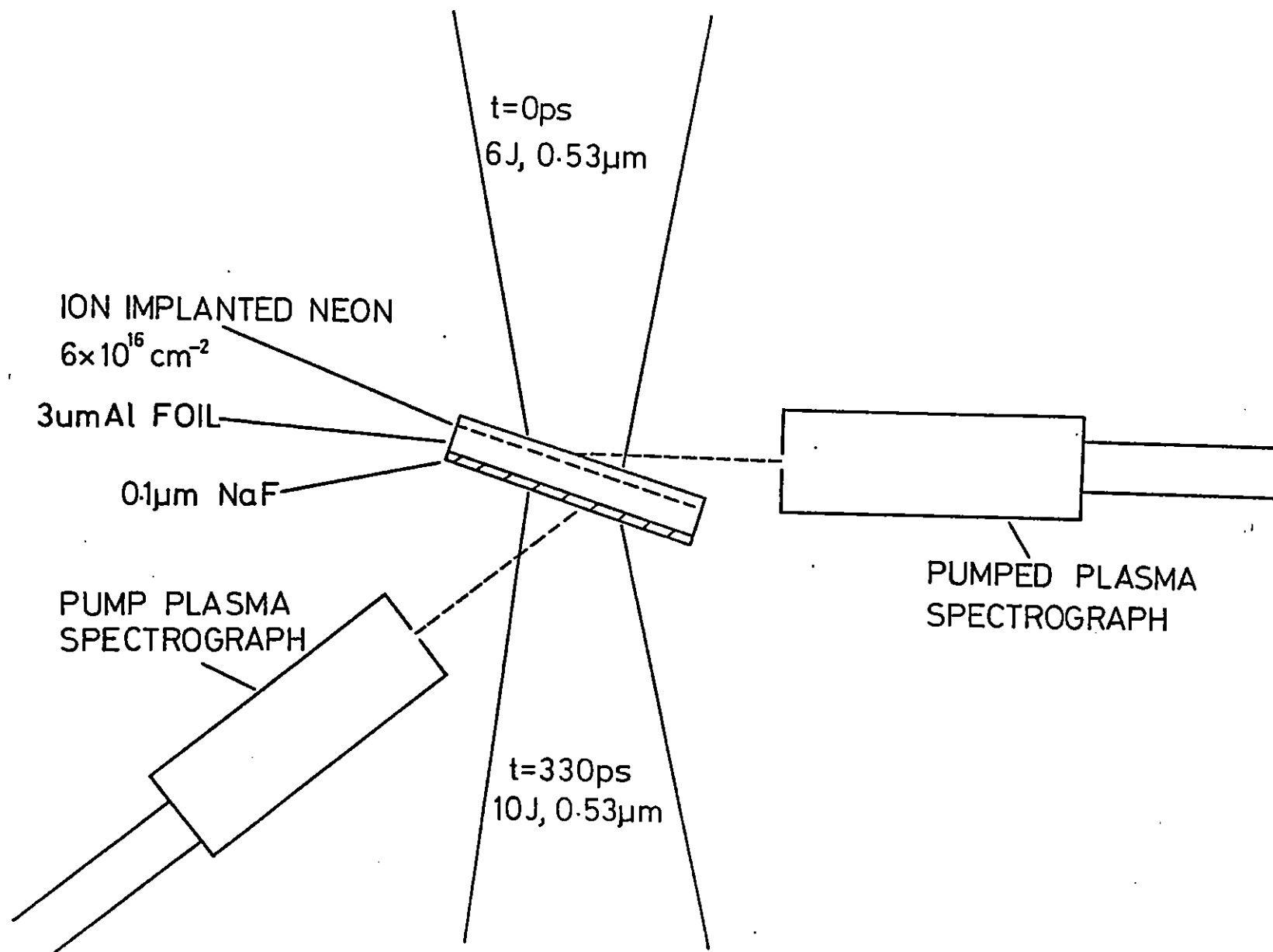


Fig. A4.18 The experimental arrangement showing the positions of the two spectrometers to measure the X-ray emission of the pump and pumped plasmas.

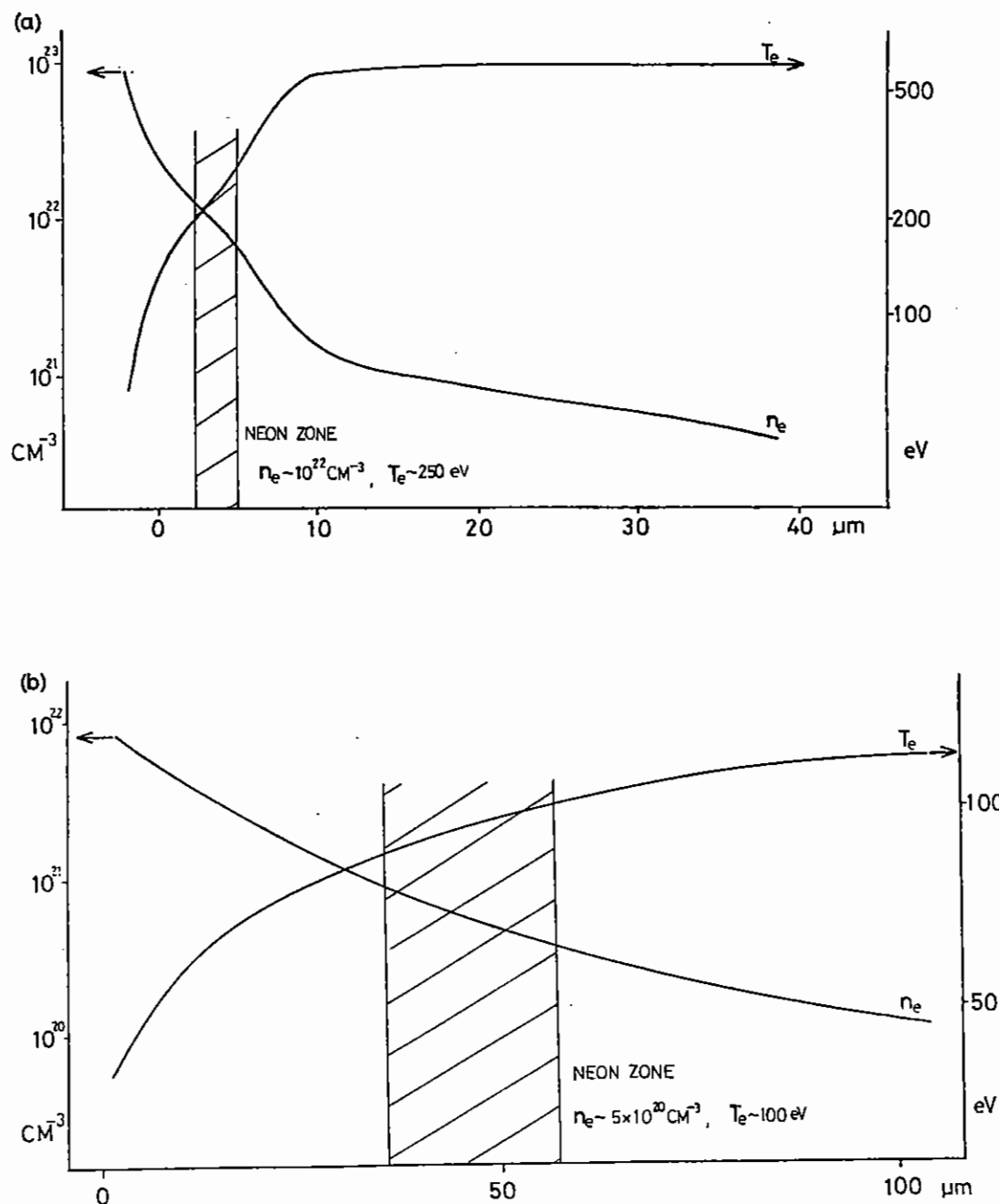


Fig. A4.19 MEDUSA calculation of the electron density (n_e) and electron temperature (T_e) at the end of the laser pulse (a), and 330 ps after the peak of the pulse (b), for a plane aluminium foil irradiated at $10^{14} \text{ W cm}^{-2}$ with a 100 ps, $0.53 \mu\text{m}$ laser pulse. The region of the foil containing the implanted neon is indicated.

The X-ray emission spectra from the neon and sodium fluoride sides of the target were recorded using two miniature spectrographs (A4.16) as shown in Figure A4.18. On the neon side a thallium acid phthalate (Tl.A.P) crystal, with $2d = 25.75 \text{ \AA}$, was used and gave a spectral coverage of $9.5 - 21.5 \text{ \AA}$. The spectrograph viewed the plasma in a direction making an angle of $\sim 10^\circ$ with the plane of the target. On the sodium fluoride side the spectrum was recorded over a similar spectral range using a rubidium acid phthalate (Rb.A.P) crystal with $2d = 26.12 \text{ \AA}$. Figure A4.20 shows a microdensitometer tracing of a single shot, space-integrated spectrum from the neon side where there was no laser pulse on the sodium fluoride side. The H-like and He-like resonance lines of neon are clearly observed in first order. The H-like and He-like resonance lines of aluminium are observed in second order and some of these overlap with some of the neon lines.

In this space-integrated spectrum the dominant contribution is due to the dense ablation phase of the plasma and so the intensity ratios of H-like and He-like lines can be used to estimate the state of ionisation and electron temperature in this phase of the plasma. From the spectrum in Figure A4.20 the intensity ratio

$$\frac{I(\text{Ne X } 1s^2 2s - 4p^2 P)}{I(\text{Ne IX } 1s^2 1s - 1s4p^1 P)} = 0.35$$

which implies that the population density ratio

$$\frac{n(\text{Ne X } 4p^2 P)}{n(\text{Ne IX } 1s4p^1 P)} = 0.49$$

Collisional-radiative calculation (A4.17) of the ionisation and excitation in a neon plasma with an electron density of 10^{22} cm^{-3} gives this population density ratio for an electron temperature of 220 eV. An electron density of 10^{22} cm^{-3} was chosen for this calculation since this is typical for the ablation plasma in short pulse irradiation (A4.16). This measured value of electron temperature for the region of the ablation plasma containing the neon, compares well with the value of 250 eV, obtained from the computer simulation at the end of the laser pulse (see Figures A4.19(a)). Referring to the ionisation calculation at $T_e = 220 \text{ eV}$, $n_e = 10^{22} \text{ cm}^{-3}$ it is found

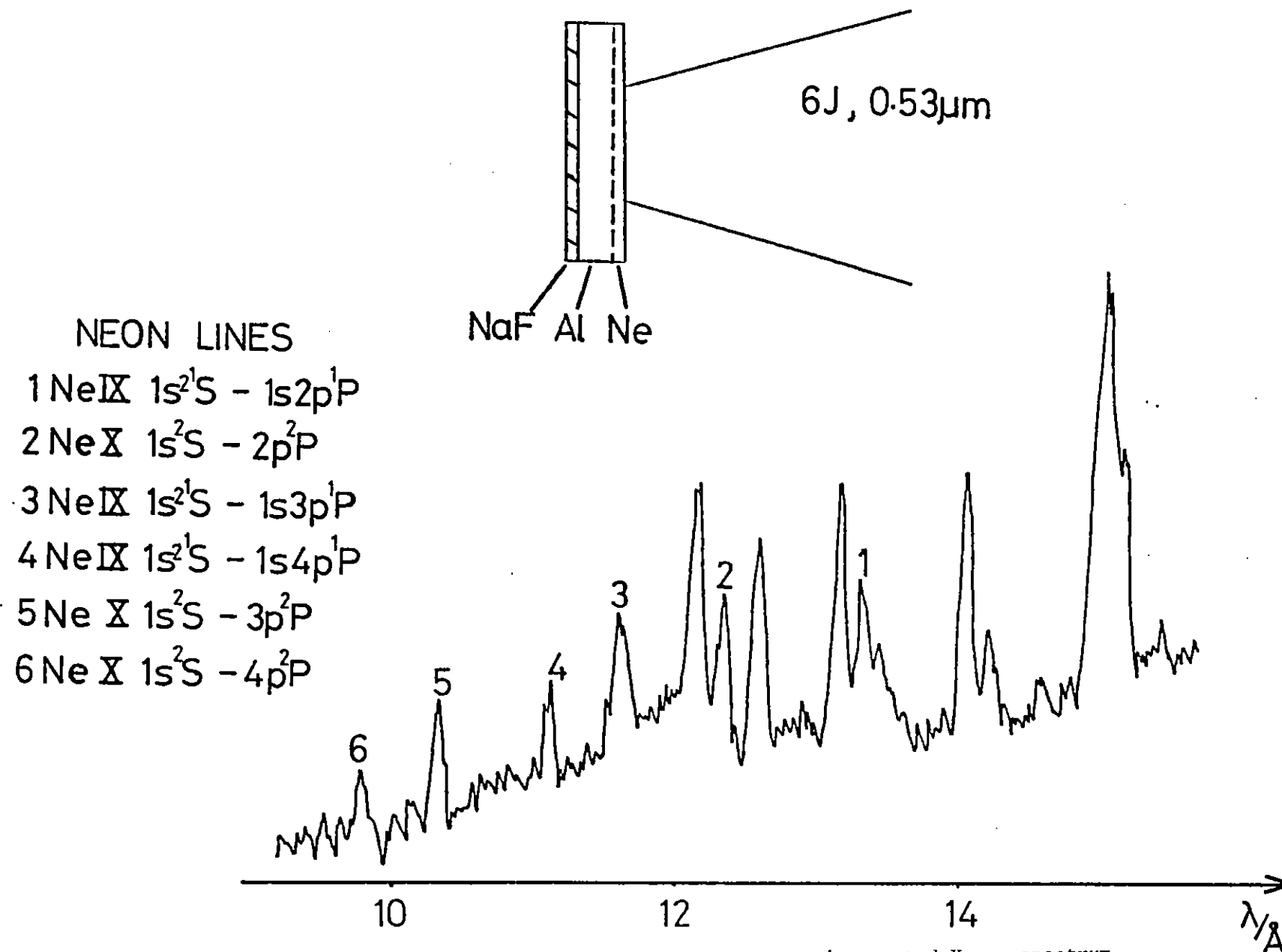


Fig. A4.20 Microdensitometer tracing of a space-integrated X-ray spectrum from the neon-implanted side of the target. The H-like and He-like neon resonance lines are indicated.

that the ion fractions are: $f(\text{Ne IX}) = 0.43$, $f(\text{Ne X}) = 0.43$ and $f(\text{Ne XI}) = 0.14$. Thus it would appear that at the irradiance used in this experiment the electron temperature is higher than the value required to maximise the fraction of He-like neon.

The absolute intensity of the Ne IX $1s^2\ ^1S - 1s4p\ ^1P$ line from the ablation plasma can be used to give an indication of the amount of neon present in the targets during irradiation. The electron temperature in the neon zone is 220 eV and at this temperature the collisional radiative ionisation model predicts that the fraction of neon ions in the $1s4p\ ^1P$ state is 4.2×10^{-4} . The MEDUSA simulation shows that the duration of the X-ray emission is ~ 80 ps. Thus taking the areal density of neon to be $6 \times 10^{16} \text{ cm}^{-2}$ the expected emission in the Ne IX $1s^2\ ^1S - 1s4p\ ^1P$ line from the 220 μm diameter spot is $1.1 \times 10^{-5} \text{ J ster}^{-1}$. This can be compared with a measured emission in this line of $6 \times 10^{-6} \text{ J ster}^{-1}$. This would indicate that a substantial fraction of the implanted neon was present when the targets were irradiated.

When the photo-pumping plasma is produced the neon-containing plasma has moved away from the initial position of the target surface. If there is measurable fluorescence on the He-like neon lines caused by the photo-pumping then there should be a change in the distribution of this line emission along the axial direction normal to the target surface when the photo-pumping is applied. X-ray spectroscopy with spatial resolution along this axis is one way of looking for evidence of photo-pumping. It was possible to record space-resolved emission spectra on the neon side by placing a 75 μm slit in the plane of dispersion of the spectrograph and at a distance of 5 mm from the plasma. Space-resolved neon spectra were recorded with and without a laser pulse on the sodium fluoride side of the target, but no clear indication of photo-pumping was observed. The pumping pulse was applied 330 ps after the laser pulse on the neon side and the computer simulation shows that at that time the neon region is at a distance of only 50 μm from the initial target surface position so that the 75 μm slit would not spatially resolve any enhanced resonance line emission that may be present in the ablation plasma emission.

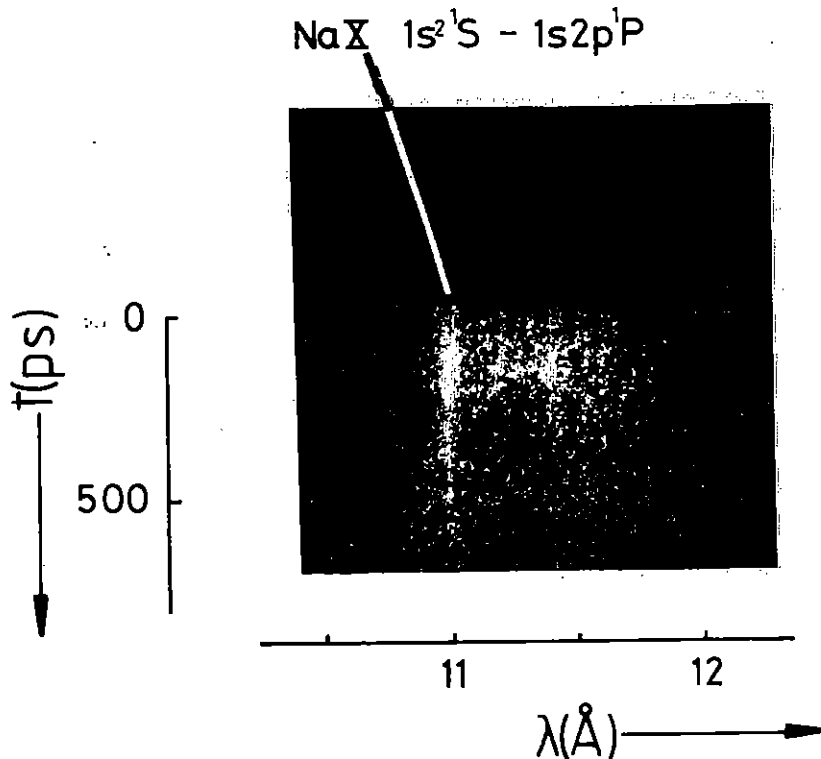


Fig. A4.21 Time-resolved and spectrally-resolved emission in the Na X $1s^2\ ^1S - 1s2p\ ^1P$ pump line from a planar Na F target irradiated at $4 \times 10^{14} \text{ W cm}^{-2}$ with a 100ps, 1.06 μm laser pulse.

A space-resolving spectrograph with a 25 μm slit was used to record the pumping spectrum from the sodium fluoride side of the target. The time integrated emission on the Na X $1s^2\ ^1S - 1s2p\ ^1P$ line was $1.2\ \text{J cm}^{-2}\ \text{ster}^{-1}$. Because of the combination of large source size and small spectrograph this line was not spectrally resolved. In a separate experiment the Na X $1s^2 - 1s2p\ ^1P$ line was time-resolved and spectrally-resolved using an X-ray streak camera at Queen's University Belfast. A similar target was used and the plasma was formed with a 100 ps 0.8 J laser pulse at $1.06\ \mu\text{m}$ focused to a $50\ \mu\text{m}$ spot giving an irradiance of $4 \times 10^{14}\ \text{W cm}^{-2}$. The time resolved spectrum is shown in Figure A4.21. This measurement of the relative spectral brightness of the Na X $1s^2\ ^1S - 1s2p\ ^1P$ line as a function of time can be used to estimate the peak absolute spectral brightness of this line in the pumping experiment. The value obtained was $6 \times 10^{11}\ \text{W cm}^{-2}\ \text{ster}^{-1}\ \text{\AA}^{-1}$ which implies a black-body temperature of 157 eV for this optically thick line. When the spectral brightness of the Na X $1s^2\ ^1S - 1s2p\ ^1P$ line is maximum the spectral half width of this opacity - broadened line is 26 m\AA, so good spectral overlap with the pumped transition is assured even when the Doppler shift due to the relative motion of the two plasmas is considered.

From a knowledge of the spectral brightness of the pump line and the fractional abundance of He-like neon in the plasma at the time of photo-pumping an estimate may be made of the expected level of fluorescence of the Ne IX $1s^2\ ^1S - 1s4p\ ^1P$ line due to photo-pumping. The time-dependent He-like neon fraction in the recombining plasma was estimated using an average-atom ionisation model to post-process the MEDUSA hydrodynamic simulation of the plasma heating and expansion (see Figure A4.22). At the time of photo-pumping, 330 ps after the peak of the heating pulse, the fraction of He-like neon was found to be ~ 15%. If the neon region of the plasma receives black-body pump radiation at 157 eV from 2π steradians then the population density ratio

$$\frac{n(\text{Ne IX } 1s4p\ ^1P)}{n(\text{Ne IX } 1s^2\ ^1S)} = 1.2 \times 10^{-3}$$

Taking the duration of the pumping pulse to be 80 ps, the expected

fluorescence in the Ne IX $1s^2\ ^1S - 1s4p\ ^1P$ line from a $220\ \mu\text{m}$ diameter spot is $5 \times 10^{-6}\ \text{J ster}^{-1}$. This is comparable to the emission recorded on this line in the ablation phase of the region containing the neon. Thus, if the line emissions from the ablation and pumped phases of the neon plasma can be spatially or temporally resolved, then it should be possible to use the fluorescence signal to monitor the pumping efficiency.

In conclusion, a novel technique of using ion-implantation of neon into a thin aluminium foil has been used to make a target suitable for the investigation of soft X-ray laser action via resonant photo-pumping. X-ray spectroscopy was used to measure the state of ionisation of the neon plasma formed by the laser irradiation of this target. Space-resolved X-ray spectroscopy was used to look for fluorescence on the Ne IX $1s^2\ ^1S - 1s4p\ ^1P$ line due to pumping by the Na X $1s^2\ ^1S - 1s2p\ ^1P$ line from a sodium fluoride plasma on the rear of the target. In this preliminary investigation pumping was not observed, but the measured spectral brightness of the pump line suggests that it should be possible to observe photo-pumping in this scheme in an optimised experiment using space or time resolution of spectral emission.

A4.3.3 Alternative Schemes

Lunney (A4.11) has proposed that transitions between certain combinations of doubly-excited levels of lithium-like ions can be used for photo-pumped soft X-ray lasers. Choosing the lower laser level to be autoionising relaxes the opacity constraint which is inherent in other photo-pumped schemes. The ion is pumped from a $1s^2\ 2l$ to a doubly-excited $1s2\ell 3p$ level which is metastable to autoionisation. The laser transition is from this $1s2\ell 3p$ level which is strongly autoionising. Taking Fluorine as an example, photo-pumping of the F VII $1s^2\ 2s\ 2s_{1/2} - 1s2s\ (^3S)\ 3p\ 2p_{2/3},\ 1/2$ transition at $14.96\ \text{\AA}$ would lead to gain at $105.6\ \text{\AA}$, the lower laser level being $1s\ 2s^2\ 2s_{1/2}$. A suitable pump line is found in the Lyman- α transition in F IX at $14.9841\ \text{\AA}$. Therefore this scheme could be investigated experimentally using the techniques described in A4.3.2 above. A suitable target would be thin mylar foil coated on each side with magnesium fluoride. Some preliminary spectra were recorded from MgF_2 coated foils irradiated on one side at $10^{14}\ \text{W cm}^{-2}$

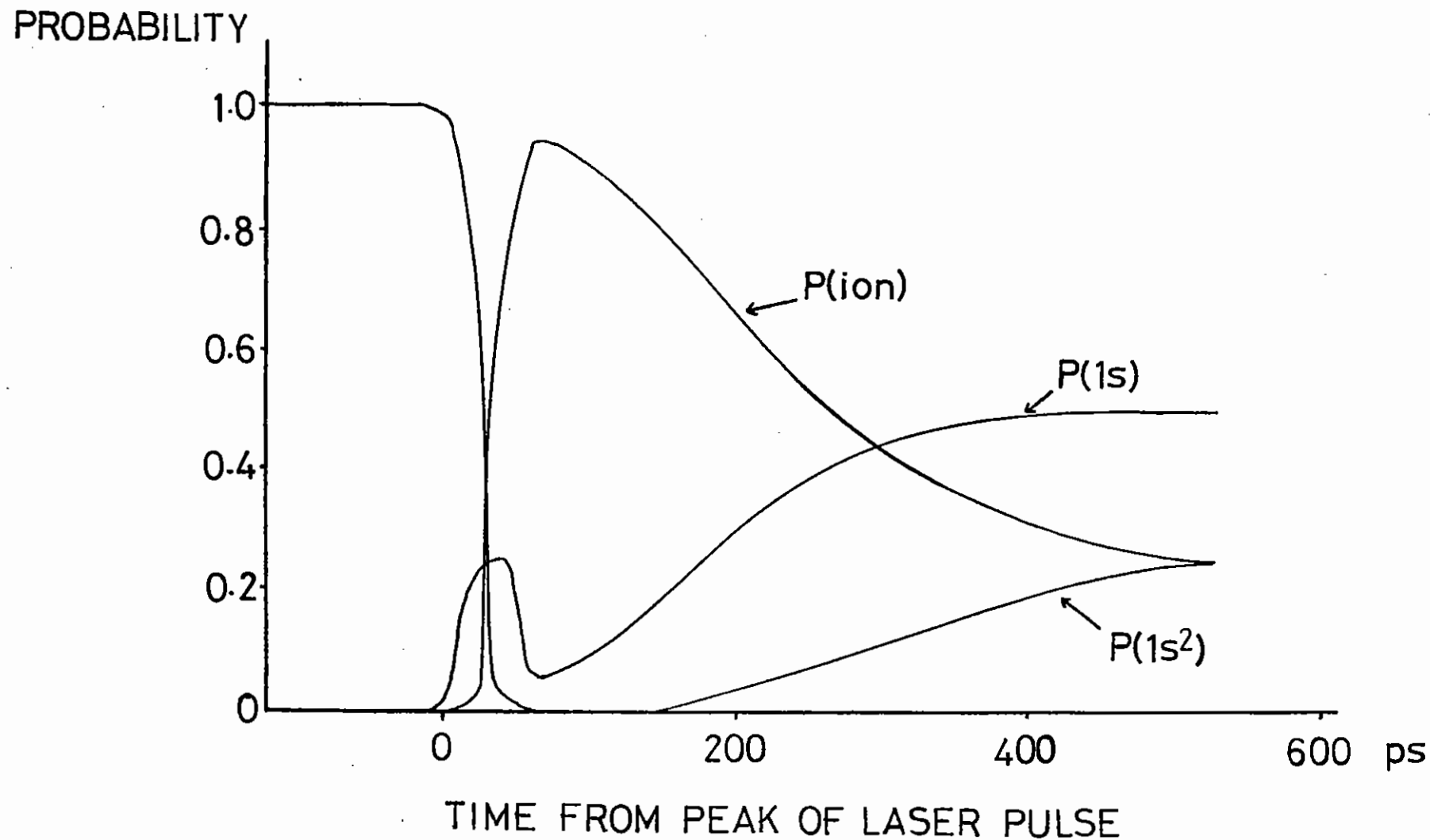


Fig. A4.22 Calculated temporal development of the fractional abundances of Ne XI ($P(\text{ion})$), Ne X ($P(1s)$) and Ne IX ($P(1s^2)$) in the neon-implanted zone of the target.

with a 100 ps, 0.53 μm laser pulse in order to determine the brightness available in the F IX Lyman- α line which was found to have an equivalent black body temperature of 130 eV.

Note that when using foil targets containing identical elements on each side then all emission lines in the pumping plasma are in coincidence, apart from a small Doppler shift, with the corresponding transitions in the pumped plasma. Therefore it should be possible to observe photo-pumping on all of these lines in the form of enhanced fluorescence, provided sufficient spatial or temporal resolution can be employed.

A common problem in all photo-pumping schemes using line sources is that of obtaining a sufficiently close wavelength match with the pumped transition. Some spectra were recorded from mylar foil targets coated with a 0.1 μm layer of chromium or tellurium in order to assess their value as alternative pump sources. Under similar irradiation conditions of $10^{14} \text{ W cm}^{-2}$ in 100 ps at 0.53 μm the black body temperatures at 15 $\text{\AA}$ were 140 eV and 110 eV for Cr and Te respectively. This is to be compared with 130 eV obtained for the F IX Lyman- α line. Although both spectra appeared as continua, that of Cr may be unresolved, closely-spaced, L-series lines (A4.18) in which case the pumping efficiency will depend on the exact values of wavelengths and linewidths. That of Te however is believed to be true continuum and so the problem of wavelength matching to the pumped transition is eliminated.

A4.4 EXAFS Spectroscopy

A4.4.1 Introduction

R W Eason (RAL)

In last years Annual Report (A4.19) we described the first EXAFS experiments carried out at the Central Laser Facility. During the past year, we have pursued the field of EXAFS in several directions. Firstly, the time integrated data on Aluminium has been analysed in conjunction with Daresbury Laboratory, using Fourier transform techniques, to yield radial structure functions, and hence derive atomic shell positions. Both cold and 'hot'

spectra have been studied.

Secondly, a further experiment has been carried out to time resolve the change previously observed in the Aluminium lattice, generated through flash heating via X-ray absorption. The new compact X-ray streak camera (A4.19) was used, and the quality of the streak pictures produced is seen to be a considerable improvement.

Lastly, a trial experiment has been performed to demonstrate the feasibility of laser reflexafs measurements. This technique utilises grazing incidence reflection from a surface rather than bulk transmission measurement to yield EXAFS spectra. The results are very encouraging and demonstrate the practicality of making measurements in this way. The fact that reflexafs is a surface-sensitive technique means that the next step is to investigate the possibility of following either slow surface dynamics (via pump and probe techniques) or fast changes (using a streak camera). It is hoped that such possibilities will attract those in semiconductor research who need to have answers to many questions about such surface effects.

A4.4.2 Radial Distribution Functions for Time Integrated Aluminium Spectra

R W Eason (RAL), D K Bradley (IC) and G N Greaves (Daresbury)

To analyse the EXAFS spectra for cold and 'hot' aluminium reported in (A4.19) the spectrum was firstly digitised into a 750 point array, and the film density converted into X-ray absorption coefficient (μ) using the film curve for SB5 film (A4.20), and published mass absorption coefficients (μ/ρ) (A4.21).

The local radial distribution of atoms can be obtained from the modulus of the complex Fourier transform of $\chi(k)$, derived from

$$\chi(E) = \frac{\mu(E) - \mu_0(E)}{\mu_0(E)}$$

and

$$E = E_0 + \frac{\hbar^2 k^2}{2m}$$

This has previously been used by Epstein and coworkers (A4.22) but places shells of atoms at distances typically 10% to 20% shorter than their true values. The reason for this misplacement is the phase change experienced by the photoelectron as it is scattered by the atomic potential firstly of the excited atom (δ) and then of the surrounding atoms (ψ). A more realistic local radial distribution has been given by Gurman and Pendry (A4.23) as $|G(r)|$, where

$$G(r) = \frac{1}{2\pi} \int_{k_{\min}}^{k_{\max}} \frac{W(k) \chi(k) \exp(-i(2kr + 2\delta + \psi)dk)}{f(k, \pi)}$$

and $f(k, \pi)$ is the backscattering amplitude. Values for f , δ and ψ for Aluminium were calculated using the program MUFPOT from the SRS program library (A4.24). $W(k)$ is a gaussian window function which for the present analysis was centred at approximately the middle of the spectral range. This is to help minimise the finite wavevector range defined by $k_{\min}$ and $k_{\max}$, which in this case was $4\text{\AA}^{-1} - 8\text{\AA}^{-1}$.

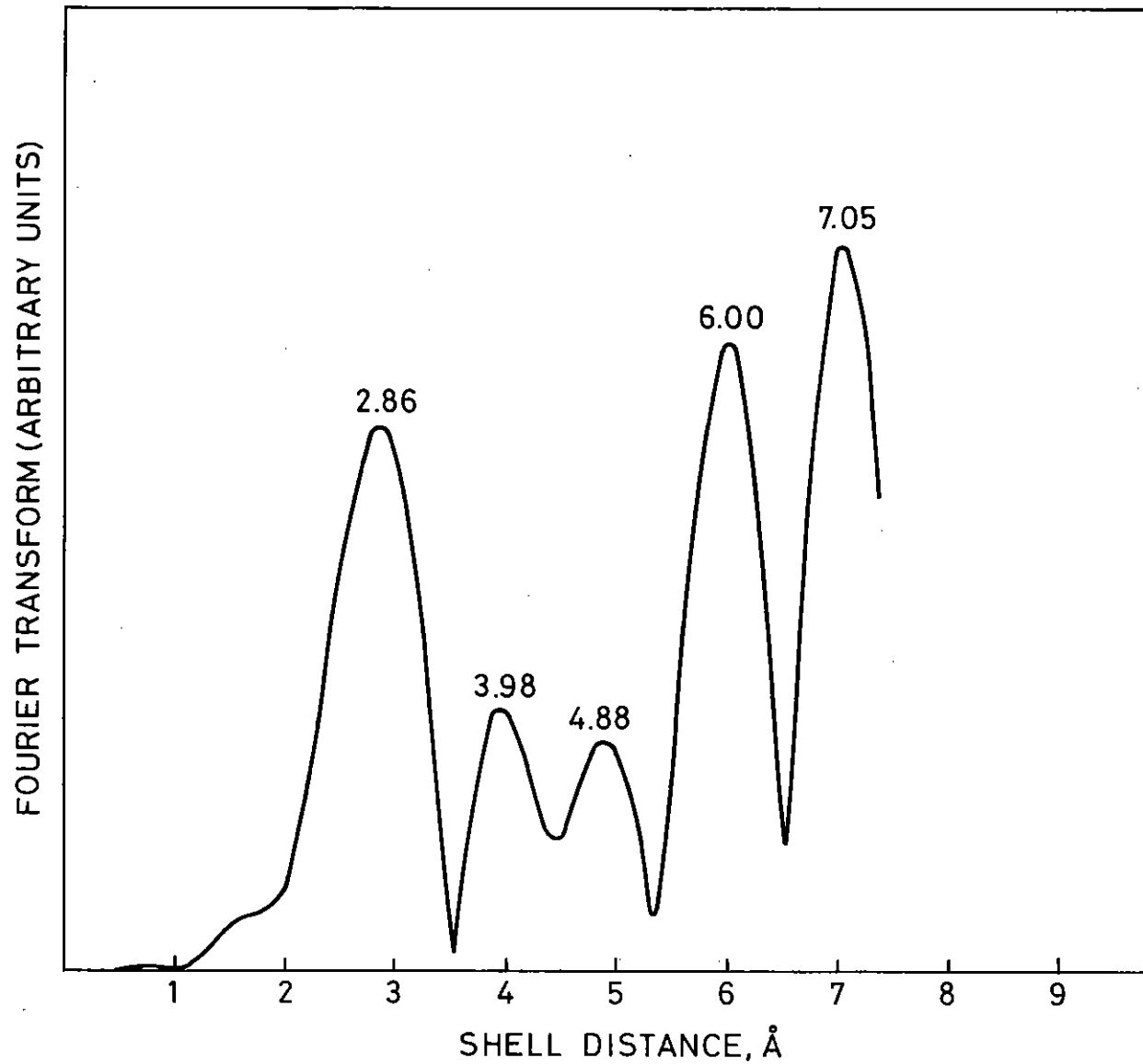
The result of the FT procedure can be seen in figure (A4.23) in which the amplitude of the FT (arbitrary units) is plotted as a function of shell distance ($\text{\AA}$). A value of 16 eV was chosen for E_0 and the precise position of $W(k)$, 170 eV, was chosen to line the first peak up with the nearest neighbour distance which for fcc Al is 2.86 $\text{\AA}$. From the table in Figure (A4.23) one can compare the experimentally determined shell positions for the other shells with the crystallographic values. The second and third shells in the laser-EXAFS distribution are located to within 1.5% of the actual distances. The fourth and fifth shells are not resolved here but their mean position closely fits the fourth peak in $G(r)$. Likewise the fifth peak at 7.05 $\text{\AA}$ lies close to the sixth shell (7.01 $\text{\AA}$) in the fcc structure. The agreement between the EXAFS radial distribution function and the crystallographic parameters is extremely good. In particular the resolution of the outer shells in Al is similar to that which can be obtained for Cu (A4.23). Other laser-EXAFS studies of Al have left even the nearest neighbour distance ambiguous.

A similar FT analysis was applied to the 'hot' spectrum, using identical parameters for the values of $k_{\min}$, $k_{\max}$, $W(k)$, etc. The result is shown in Figure (A4.24). This radial distribution function looks somewhat different, as now there are only three peaks. The suggestion is that a dilation has occurred in the aluminium lattice, shifting the peaks to greater distances. The ratio of the first peak distance in the 'hot' case to the first peak distance for cold aluminium is 1.22. The ratio for the second peak is 1.20, a very similar value. If a comparison is made between the third peak at 6.64 $\text{\AA}$, and the third peak in the cold case, correlation is not so good. A value of 1.22 is obtained for the ratio however, if 6.64 $\text{\AA}$ is compared with the average of peaks three and four in Figure (A4.23).

The view that a rapid dilation has occurred is supported by an approximate calculation based on the kinetic energy of Aluminium atoms which have rapidly received sufficient energy to rapidly change their phase from solid to gaseous. At the boiling point of Aluminium, the average velocity of the atoms is 0.5 μm in 1 ns. If a foil of aluminium $\sim 2 \mu\text{m}$ thick were converted to a gaseous state via flash thermal heating the outside atoms would freely stream away, while the atoms in the middle of the foil would suffer more constraint. An average spacing after $\sim 1\text{ns}$ therefore of the atomic distribution would be very close to the measured value of 1.2.

XANES (X-ray absorption near edge structure) calculations on a dilated lattice also produces a feature which has been observed experimentally. A strong 'white line', a sharp dip in the value of the absorption coefficient just after the absorption edge, appears in the XANES model. This is seen in the 'hot' aluminium EXAFS of Figure 7.52 of last year's Annual Report (A4.19). It would seem therefore that the EXAFS and the XANES are in good agreement. Both the EXAFS results for cold aluminium (A4.25) and the XANES modelling for 'hot' aluminium (A4.26) will shortly be published.

FOURIER TRANSFORM OF NORMALISED EXAFS SPECTRUM OF ALUMINIUM



SHELL	RADIUS (Å)	
	THEORY	EXP.
1	2.86 ($\sqrt{a}$)	2.86
2	4.04 ($\sqrt{2a}$)	3.98
3	4.95 ($\sqrt{3a}$)	4.88
4	5.72 ($\sqrt{2a}$)	6.00
5	6.40 ($\sqrt{5a}$)	
6	7.01 ($\sqrt{6a}$)	7.05

a = NEAREST NEIGHBOUR
DISTANCE

Fig. A4.23 Fourier transform of normalised EXAFS spectrum of aluminium.

FOURIER TRANSFORM OF NORMALISED EXAFS SPECTRUM OF HEATED ALUMINIUM

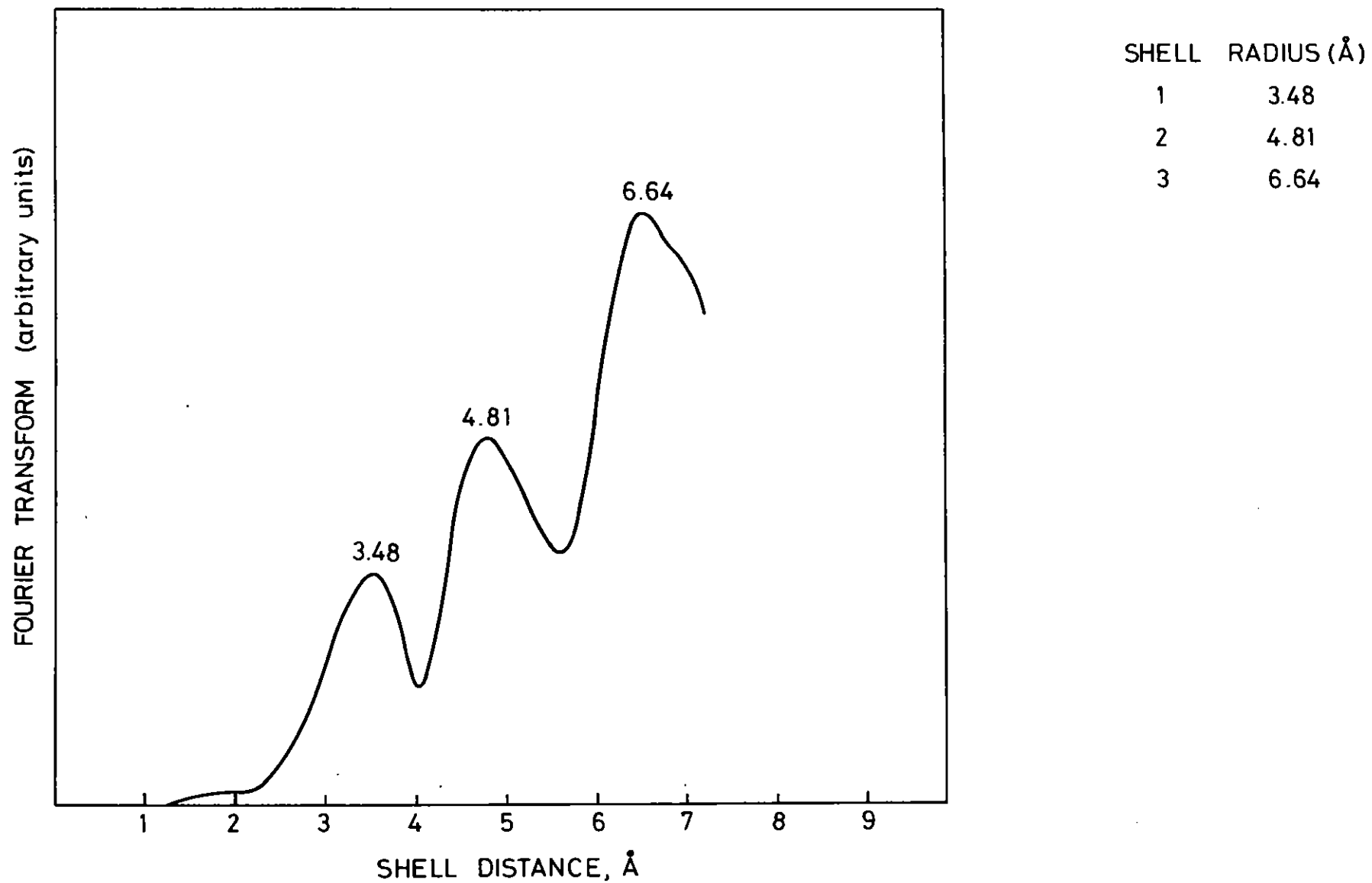


Fig. A4.24 Fourier transform of normalised EXAFS spectrum of heated aluminium.

A4.4.3 Time resolved EXAFS

R W Eason (RAL) and D K Bradley (IC)

The experiments reported in last year's Annual Report to time resolve the change that occurs during the transient heating of aluminium foil have been repeated this year, using the new compact streak camera. Problems occurred in the analysis of the data taken in an earlier experiment as numerous defects on the streak pictures, due to instrumental factors in the streak camera/intensifier combination, prevented a proper quantitative analysis. The data taken in the latest experiment is free from these artifacts, and will be analysed as soon as possible.

The experimental arrangement is virtually identical to that reported in section 7.9.3 of last year's Annual Report. The original intention had been to look at numerous samples, such as Aluminium Chloride, which undergoes a dramatic 83% volume change on melting, going from $AlCl_3$ to Al_2Cl_6 , and Lithium Bromide, looking at the Bromine L -edge, but all the samples of interest defied manufacture into suitable thin films.

The data that was taken on Aluminium foil is shown in Figures (A4.25) to (A4.27). These show the change that occurs in the EXAFS due to radiative heating, as the foil is moved progressively nearer to the plasma.

At 25 mm away from the plasma, the foil experiences little or no heating. Shot 05150384, seen in Figure (A4.25) shows a cold sample in which clear EXAFS is seen, lasting as long as the pulse stays on. A microdensitometer scan clearly shows the characteristic Al EXAFS spectrum. The fiducial is a crystal artefact, and provides a convenient wavelength calibration. This is necessary in some cases, as in this shot for example, where there is distortion in the streak camera tube due to spacecharge effects, causing the normally straight K-edge to 'warp' at high X-ray fluxes.

At 4 mm distant from the plasma, the foil will experience considerable heating. Figure (A4.26) shows shot number 14270884 in which there is only evidence of change perhaps in the last scan at $t + 600$ ps. When the foil is

moved to 3 mm away, as in shot 16270384, Figure (A4.27), then the EXAFS appears changed from the very start of the pulse. These last two figures are of shots identical in all respects apart from the plasma-to-foil distance. All three are due to be analysed using the EXAFS codes, at Daresbury. Time integrated spectra and plasma calorimeter data taken on these, and similar, shots will also provide quantitative information on the rates of heating of the foil.

A4.4.4 Laser-Reflexafs

R W Eason (RAL), D K Bradley, J D Hares and P J Dobson (IC)

A4.4.4(a) Introduction

Trial experiments have recently been carried out to investigate the feasibility of measuring reflection EXAFS spectra using a laser-produced-plasma as a source of X-rays. The technique of reflection EXAFS, or reflexafs, consists of measuring X-rays which have been specularly reflected from a surface at small angles of incidence ($\sim 1^\circ$). The reflectivity is high ($> 50\%$) below a critical angle defined by

$$\theta_c = (2\delta)^{1/2}$$

where $1-\delta$ is the real part of the refractive index, μ . Below this critical angle, the penetration depth is small (≈ 20 Å) (A4.27), while for angles greater than the critical angle, the penetration depth increases substantially to the reciprocal of the absorption coefficient, and the reflectivity falls rapidly to zero. For angles near the critical angle, fine structure may be observed in the spectrum of reflected X-rays above an absorption edge, yielding EXAFS which is similar to that obtained in transmission mode.

The reflexafs technique has the advantage that the small penetration depth involved makes it very sensitive to surface structure. Numerous surface phenomena are therefore open to investigation by this method, and the

SHOT 05150384

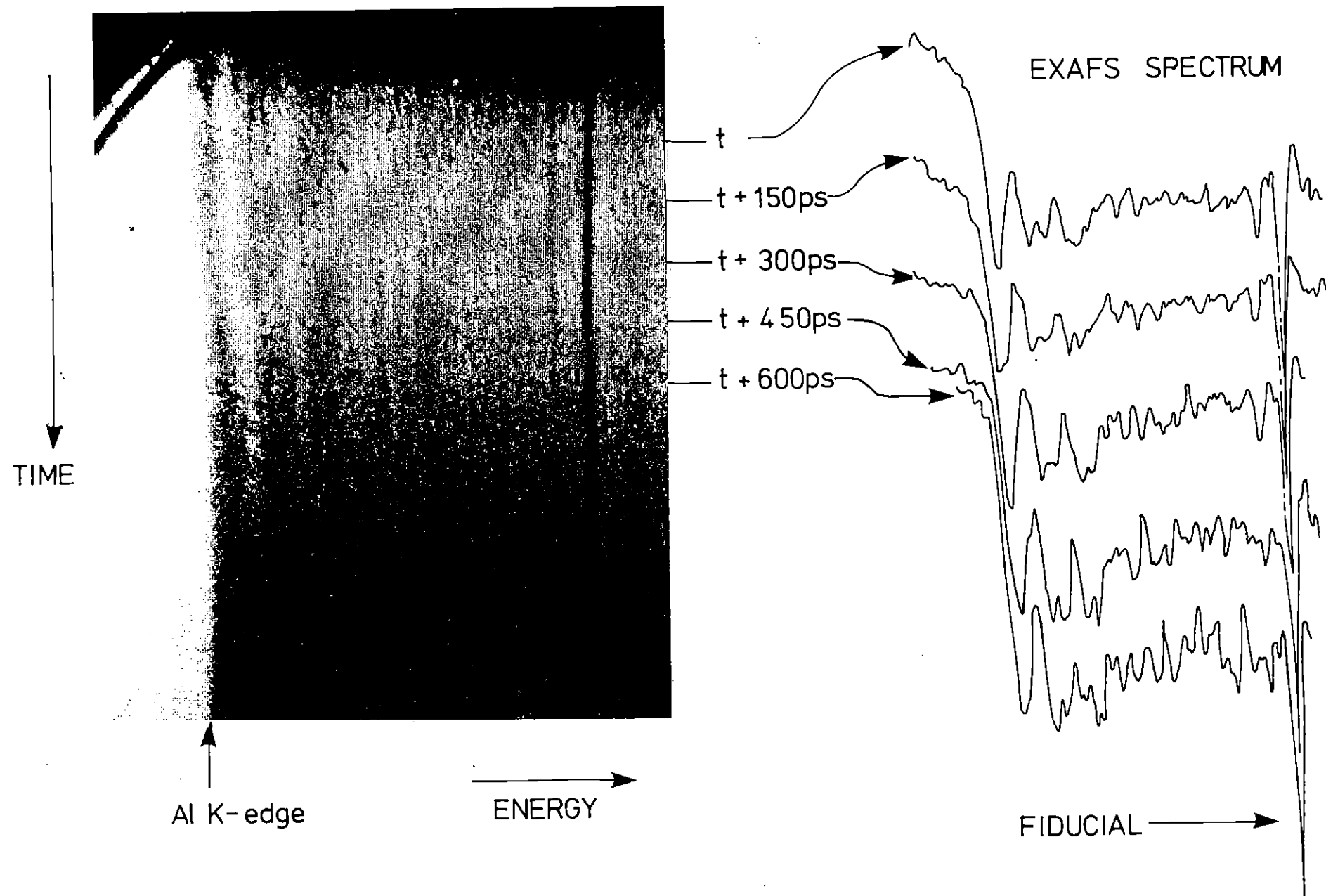


Fig. A4.25 Time resolved EXAFS spectrum of Al sample 25mm from plasma.

SHOT 14270384

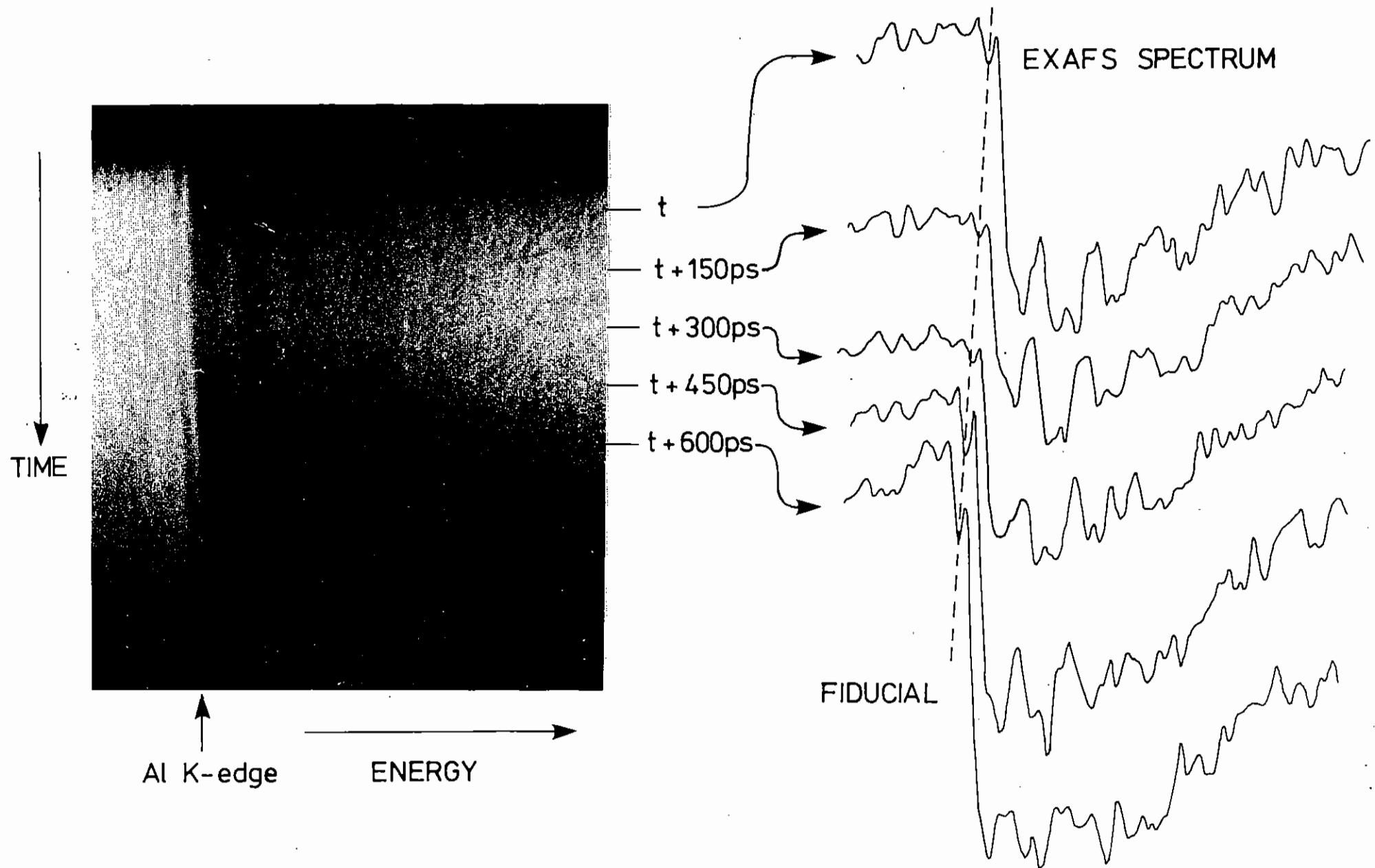


Fig. A4.26 Time resolved EXAFS spectrum of Al sample 4mm from plasma.

SHOT 16270384

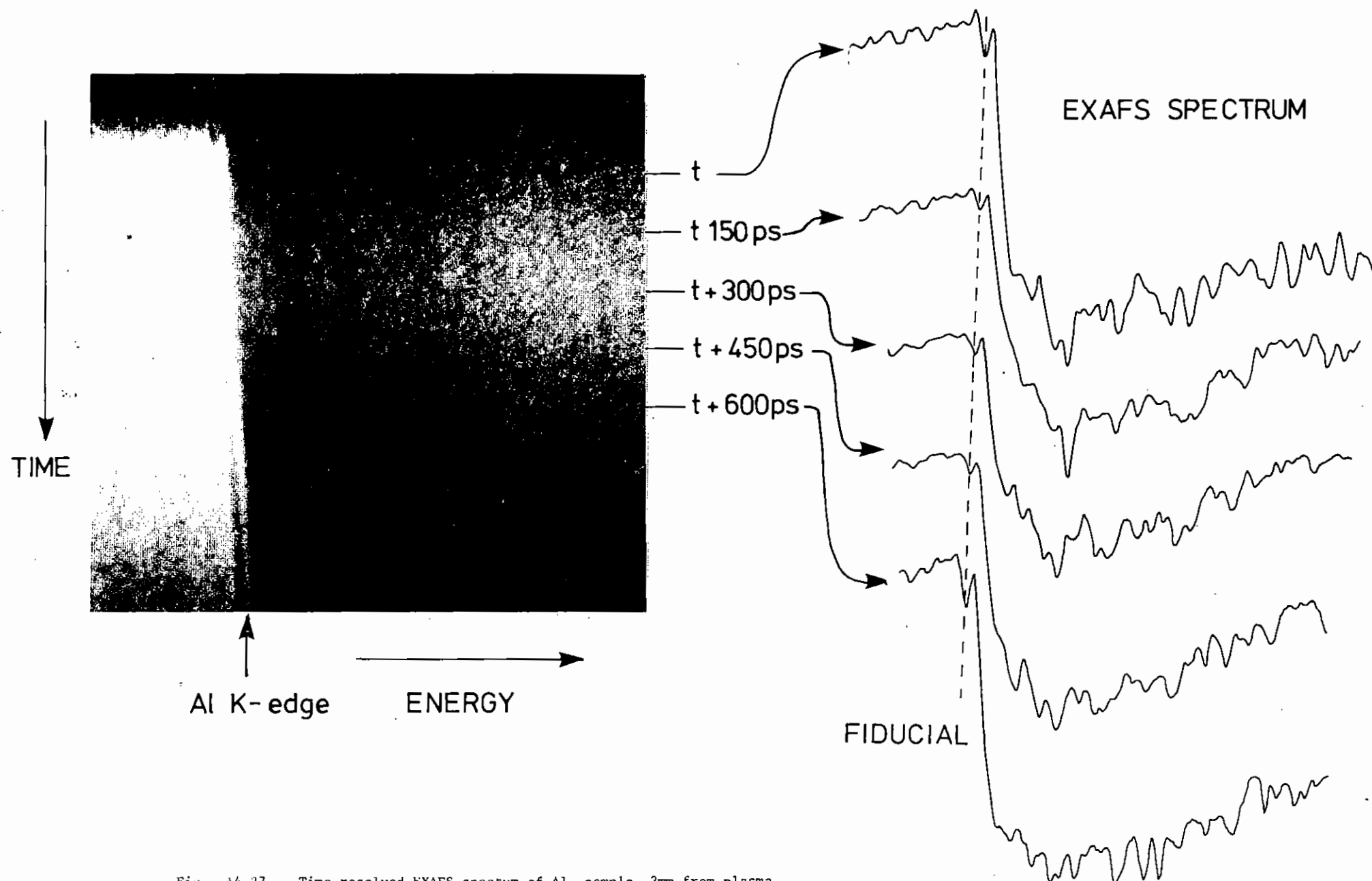


Fig. A4.27 Time resolved EXAFS spectrum of Al sample 3mm from plasma.

extremely good sample access that is inherent to this technique, makes reflexafs very attractive compared to transmission mode EXAFS for time resolving surface dynamics, or fast phase changes in surface films.

A.4.4.4(b) Experimental

The experimental arrangement is shown in figure (A4.28). A silicon wafer, whose surface roughness was flat to better than $0.05\text{ }\mu\text{m}$ over all its surface area was used as a mirror surface for grazing incidence reflection. For recording Aluminium reflexafs, a $\sim 1\text{ }\mu\text{m}$ overcoating was vacuum deposited onto the silicon surface; Si reflexafs spectra were also recorded from an uncoated wafer. Care was taken to ensure the wafer was fixed to the microscope slide in such a way not to bend it or stress it. The bismuth target used was formed into a wedge shape and fixed to the slide adjacent to the wafer. The wedge shape was used so that specific angles of incidence could effectively be chosen by initially focusing the alignment laser beam onto the microscope slide, defocusing the lens by a known amount, and then traversing the whole slide until the target came into sharp focus. This procedure allowed the height of the target above the silicon wafer to be accurately controlled. In practice an angle of incidence of 1° , and target to mid-point of wafer distance of 35 mm, meant the target was shot $\sim 600\text{ }\mu\text{m}$ above the wafer surface.

A minispectrometer equipped with an ADP ($2d = 10.65\text{ }\text{\AA}$) crystal allowed both the Aluminium and Silicon K-edges to be viewed. From the geometry involved, the film obtained showed wavelength (energy) in one dimension and angle of incidence and hence spatial position on the wafer in the other. A direct X-ray block was positioned at various heights above the silicon wafer, varying from $100\text{ }\mu\text{m}$ to $300\text{ }\mu\text{m}$.

40-50 J, 1 ns. $0.53\text{ }\mu\text{m}$ laser pulses were used, focused down to $\sim 200\text{ }\mu\text{m}$ spot size, to obtain high brightness X-ray sources. The film used was Kodak DEF, which had been observed to have a very low background fog level.

The results are shown in Figure (A4.29). In (a) the direct X-rays can be seen as well as the dispersed spectrum. The curved K-edge that can be seen

in both (a) and (b) is real, and is the result of anomalous dispersion (δ/λ^2 varies with λ). This effect has been seen before, and appears in the data taken using synchrotron sources (eg (A4.28)). It is not so immediately obvious in other reflexafs work however, as most other spectra have been recorded in a sequential manner, ie at one angular position at a time, whereas the laser-reflexafs technique is best suited to sampling a range of angles in a single shot.

For $\lambda \ll \lambda_k$ and $\lambda \gg \lambda_k$ the value of δ/λ^2 is $\sim$ constant, and so θ_c is proportional to λ . As λ approaches the K-edge however, δ/λ^2 starts to change rapidly. The value of δ has been calculated as a function of λ using both the Kramers-Kallmann-Mark theory, and the quantum dispersion theory of Hönig (A4.29), and has been used to evaluate the critical angle, θ_c as a function of energy for Aluminium at $\lambda = \lambda_k$. The results are shown in figure (A4.30) in which both axes have scales which are the same as the experimentally obtained values for the spectra in Figure (A4.29).

If the two spectra (a) and (b) of figure (A4.29) are overlaid onto the two theoretical curves, then it is seen that the theory of Hönig and the experimental data are in close agreement. This comparison allows accurate angles of incidence, and hence critical angles, θ_c , to be assigned to the two spectra. It can also be seen that (a) has been cut off on both sides, ie at both low and higher angles of incidence, whereas (b) has only suffered cut-off on the low angle side. This cut-off is due to the exact positioning of the silicon wafer, and the direct X-ray block.

The spectra can be compared with other EXAFS spectra taken in the conventional transmission mode. In Figure (A4.31) three microdensitometer traces (a), (b) and (c) have been made of shot number 02290384. (a) was taken over half the width of the spectrum, from the mid-point to the higher angle of incidence side. (b) was taken over the other half, from mid-point to the low angle cut-off side, while (c) is over the whole spectrum width. The apparently decreasing abruptness of the K-edge in (a) is clearly seen, when compared to (b). Trace (d) is taken from (A4.30) which was recorded with a conventional X-ray set using a 20 hour exposure. The spectrum (d) has been re-scaled in both energy and absorption to allow direct comparison

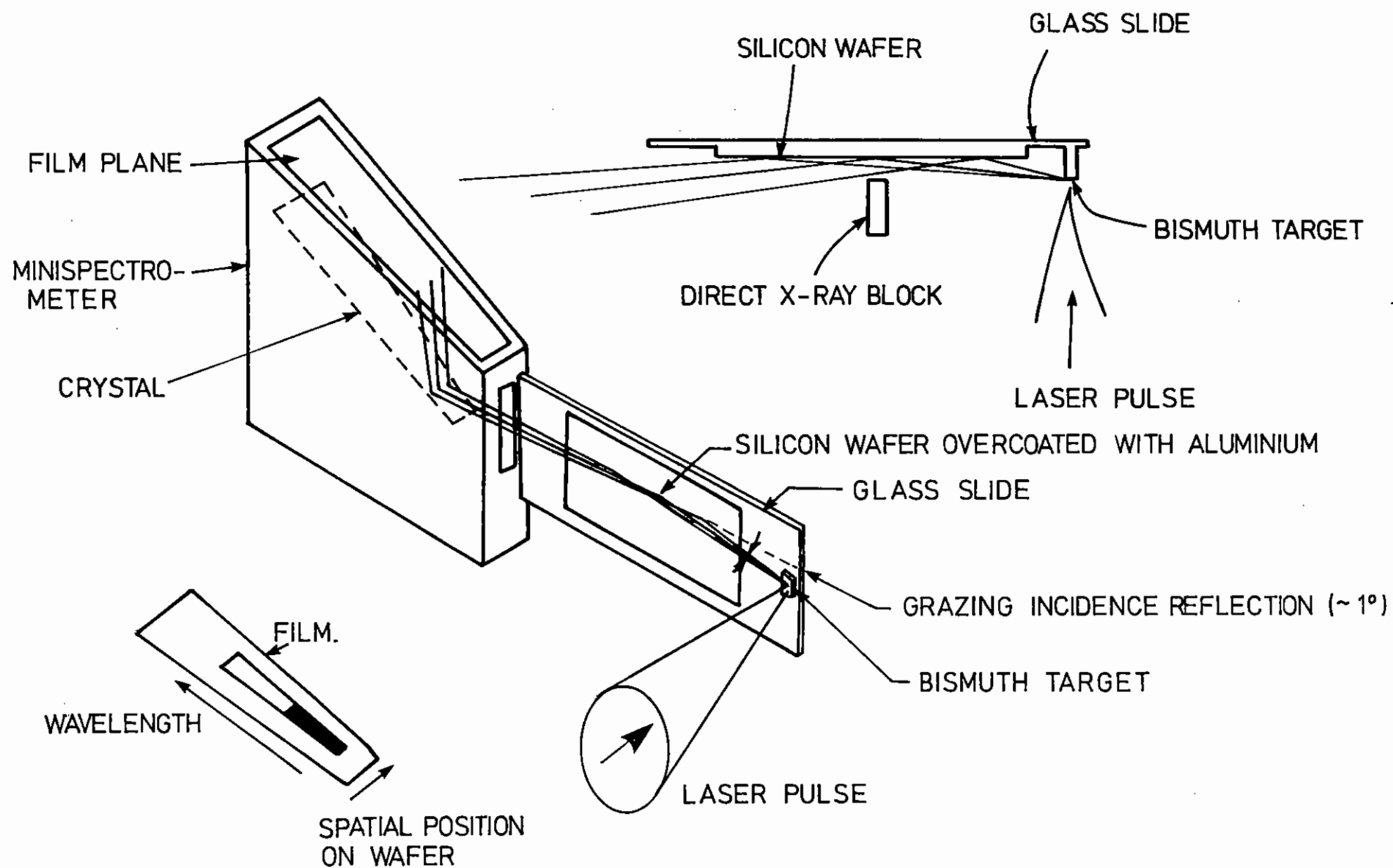


Fig. A4.28 Experimental arrangement for recording laser-reflexafs spectra.

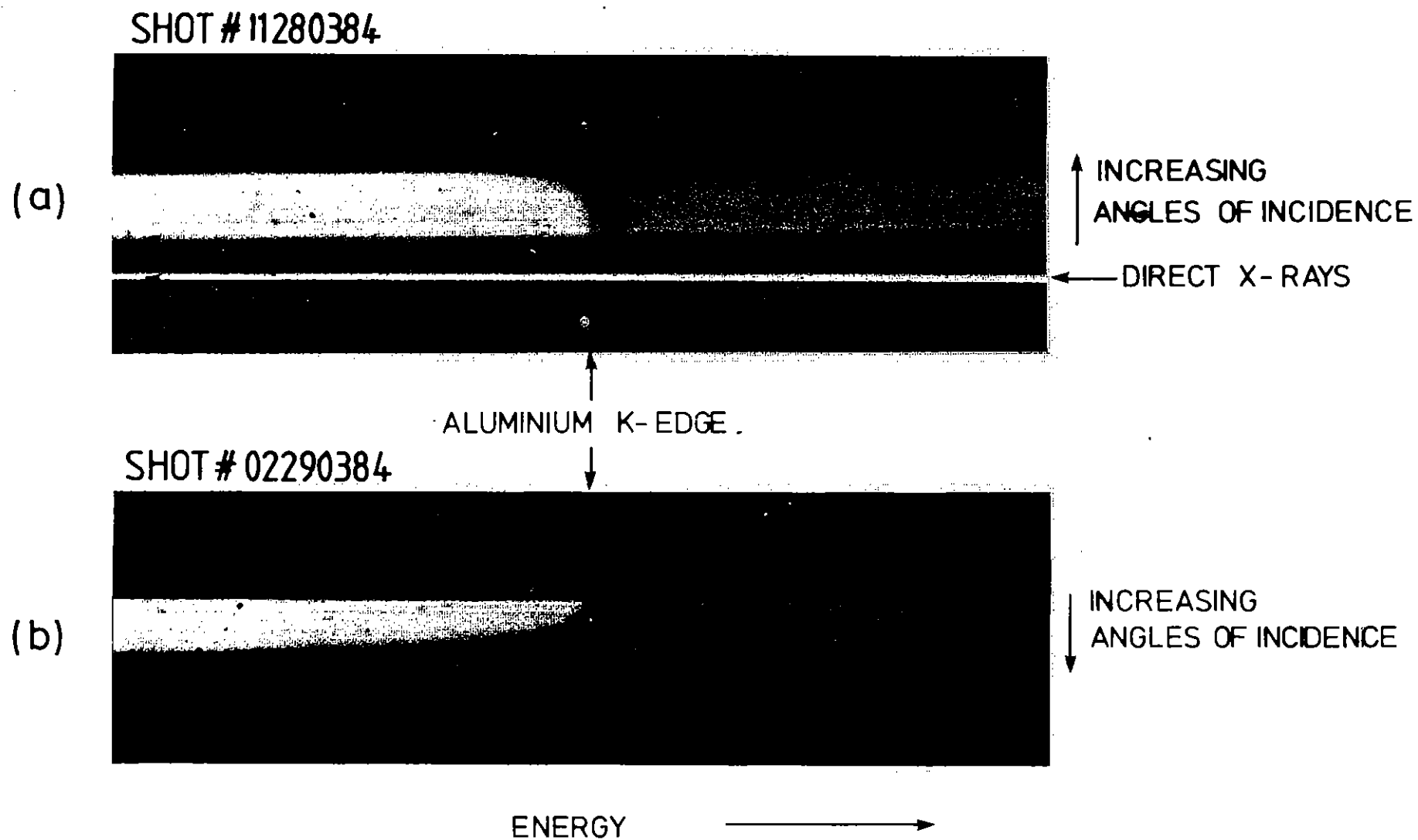


Fig. A4.29 Aluminium K-edge reflexafs spectra.

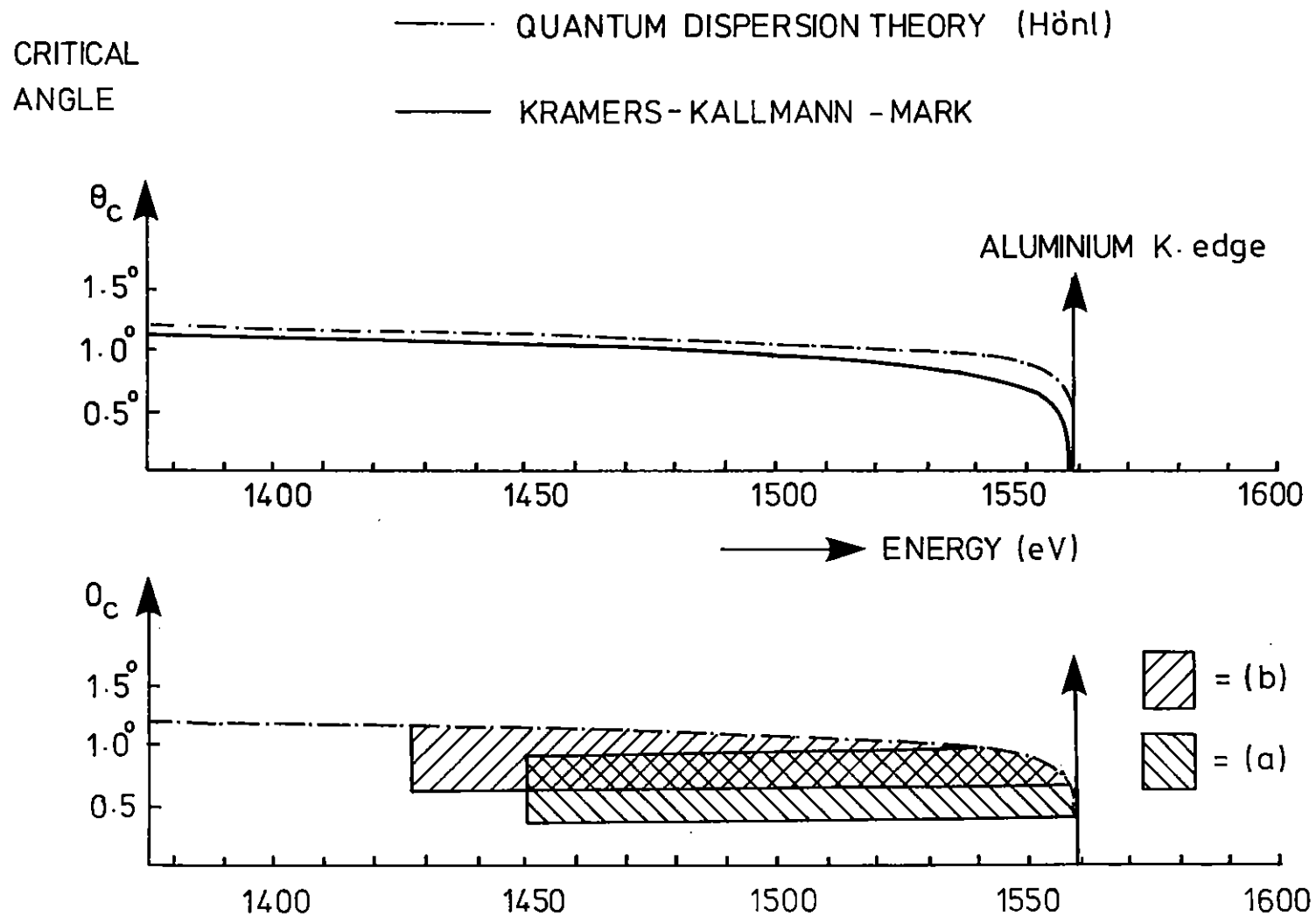


Fig. A4.30 Critical angle for reflection, θ_c , plotted as a function of energy for $E < E_K$.

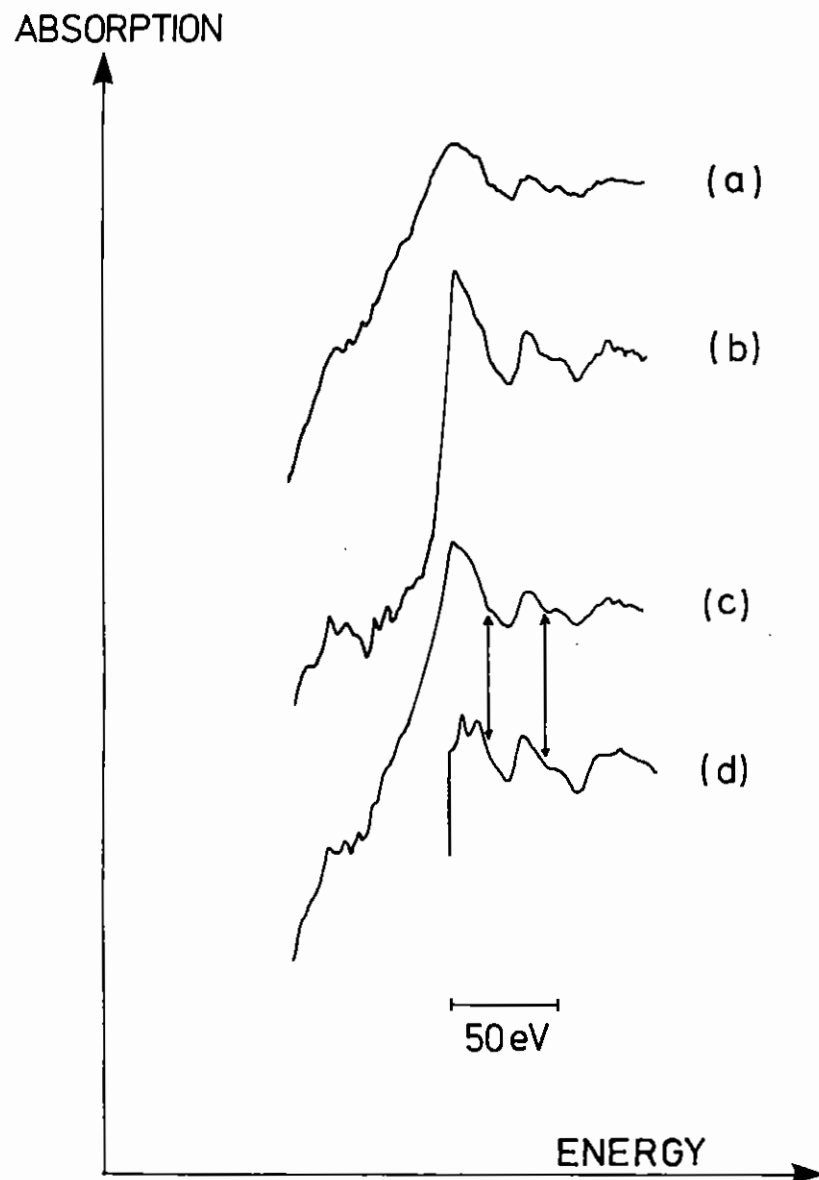


Fig. A4.31 Aluminium reflexafs spectra (a)-(c) of shot 02290384, compared to transmission mode EXAFS spectrum (d).

with (c). It can be seen that there is very good agreement between (c) and (d). The small features (arrowed) due to more distant shells of atoms in the aluminium lattice are well resolved, although the XANES structure in (d) is clearly better than (c).

This close agreement, demonstrating the ability to record good reflexafs spectra using a laser-produced plasma source is very encouraging. Further experiments are planned using either time resolved or pump-and-probe techniques.

A4.5 Soft X-ray Microscopy

R Rosser and K Baldwin (IC)

A4.5.1 Introduction

Electron microscopy of biological material suffers from the severe disadvantage of giving very low contrast between protein and water. Only dehydrated specimens are viewed, which is a problem in life sciences as ~ 80% of biological material is water and removing it conceivably results in gross structural deformation (A4.31).

In the soft X-ray region, between the oxygen absorption edges at 23 Å and carbon edge at 44 Å, protein shows marked contrast in water. Much work is being done to develop soft X-ray microscopy using synchrotron radiation as the source and zone plates as focusing elements. Two prototype microscopes now exist, one at the BESSY storage ring in Berlin (A4.32) and one at the NSLS VUV ring at the Brookhaven National Laboratory, New York (A4.33). These microscopes are anticipated to have an ultimate resolution of 100 Å, limited by radiation damage to the specimen. By taking the image in a very short time, ~ 1 nsec, it may be possible to take the picture before the damage has time to manifest itself. The resolution would then be limited only by the optics and could perhaps be made as small as 20 Å (A4.34).

To look into this possibility of using a short-lived laser produced plasma as an X-ray source for soft X-ray microscopy, an experiment was set up in

the double beam target area of the RAL glass laser facility in August 1983.

A4.5.2 Experiment

A 1 nsec 0.53 μm , 40 J pulse was focussed down onto a planar carbon target. The CV (40.3 Å) line of the resultant carbon plasma was used to illuminate a specimen. A 36 mm focal length zone plate was used to produce an image of the specimen on Kodak 101-01 film, in the shadow of a gold wire, as illustrated in Figure A4.32.

The system was aligned and spectra (Figure A4.33) indicated sufficient light for image formation with the CV line. However no image was produced (Figure A4.34).

Subsequent examination of the zone plates at the Brookhaven National Laboratory, revealed a 5-10 μm contamination layer on the surface of the zone plate. We believe the absorption of soft X-rays by this layer to be the reason for the failure of the imaging. It is hoped that it will be possible to repeat the experiment, without contamination, in 1984.

A4.6 Pulsed Laser Plasma X-ray Diffraction Study of crystallographic Changes During Laser Annealing

P J Dobson, J D Hares, S D Tabatabaei (IC), R W Eason (RAL) and J G Lunney (Dublin)

There is currently a great deal of interest in the transient structural effects that occur during laser annealing of materials, and in particular, semiconductors. This report describes the results of a feasibility study to investigate changes that occur in the near surface region of silicon during laser annealing. Previous structural studies in this area are restricted to work by Larson et al (A4.35) (Oak Ridge Laboratory) using a synchrotron source of X-rays and a time-gating of the laser annealing pulse and the synchrotron X-ray pulse. The data produced by that group represents an important landmark in the field, providing for the first time a description of the temperature distribution with depth for various times after the heat

pulse. The results are also considered as rather surprising by many workers (eg A4.36), particularly with regard to the cooling rates which were observed.

In the experiments at RAL a simple X-ray diffraction apparatus as shown in Figure (A4.35) was mounted in the target chamber. The main laser pulse was split into two beams, one of which gave a heating pulse of $\sim 1 \text{ J cm}^{-2}$. This heating beam was homogenised by a subsidiary plasma by focusing with an f/20 lens onto a mylar target at 5°. The specularly reflected light, having lost any coherence because of plasma scattering was relayed onto the Si sample with another lens. While this removed any fine scale speckle the overall uniforming however was not better than 30%. The other beam was time delayed and then focused onto a suitable target to produce quasi-monochromatic X-rays. For most experiments a chlorine He α -like complex of lines $\sim \lambda = 4.44 \text{ Å}$ was used. A portion of the silicon crystal specimens was shielded from the heating beam, but still sampled with the X-ray beam. This was in order to provide an unperturbed set of diffraction features on the X-ray film. Typical results can be seen in Figures A4.36 and A4.37. These results show an easily measurable shift of the diffraction angle and a broadening of the features. The shift and broadening were at a maximum for times of between 20 to 80 ns after the heating beam. The shift of the diffracted angle corresponds to an increase in the lattice parameter up to around 0.7%. The broadening has several origins, viz: contributions of X-rays scattered from different depths, corresponding to different temperatures and hence lattice spacings; thermal diffuse scattering from the hot lattice; scattering from the liquid surface layer. We cannot at present be very specific about these contributions, but we have sufficient data and information to make a more detailed analysis.

The advantages that our experimental arrangement has over the previous study by Larson et al (A4.35) are as follows:

- (a) All of the information is obtained in a single laser shot, ie we do not require averaging over several thousand shots.
- (b) The wavelength chosen by us is much more appropriate for determining the

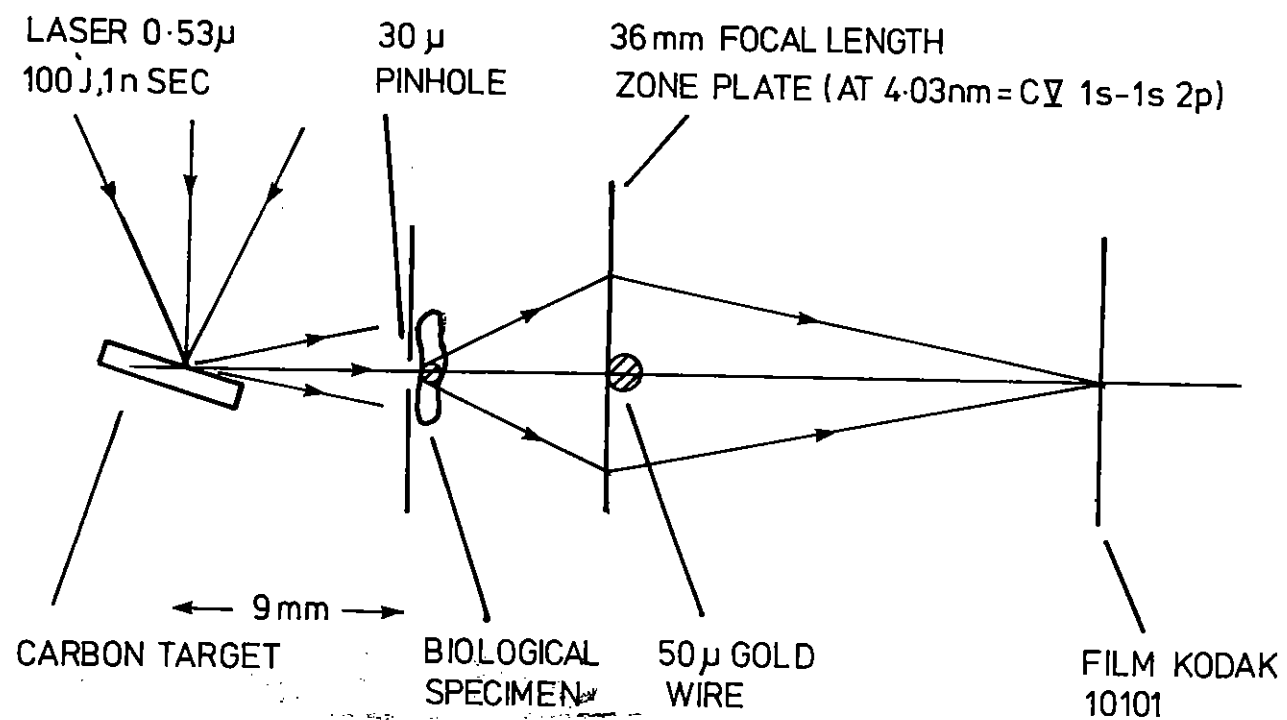


Fig. A4.32 Experimental arrangement used for soft X-ray imaging.

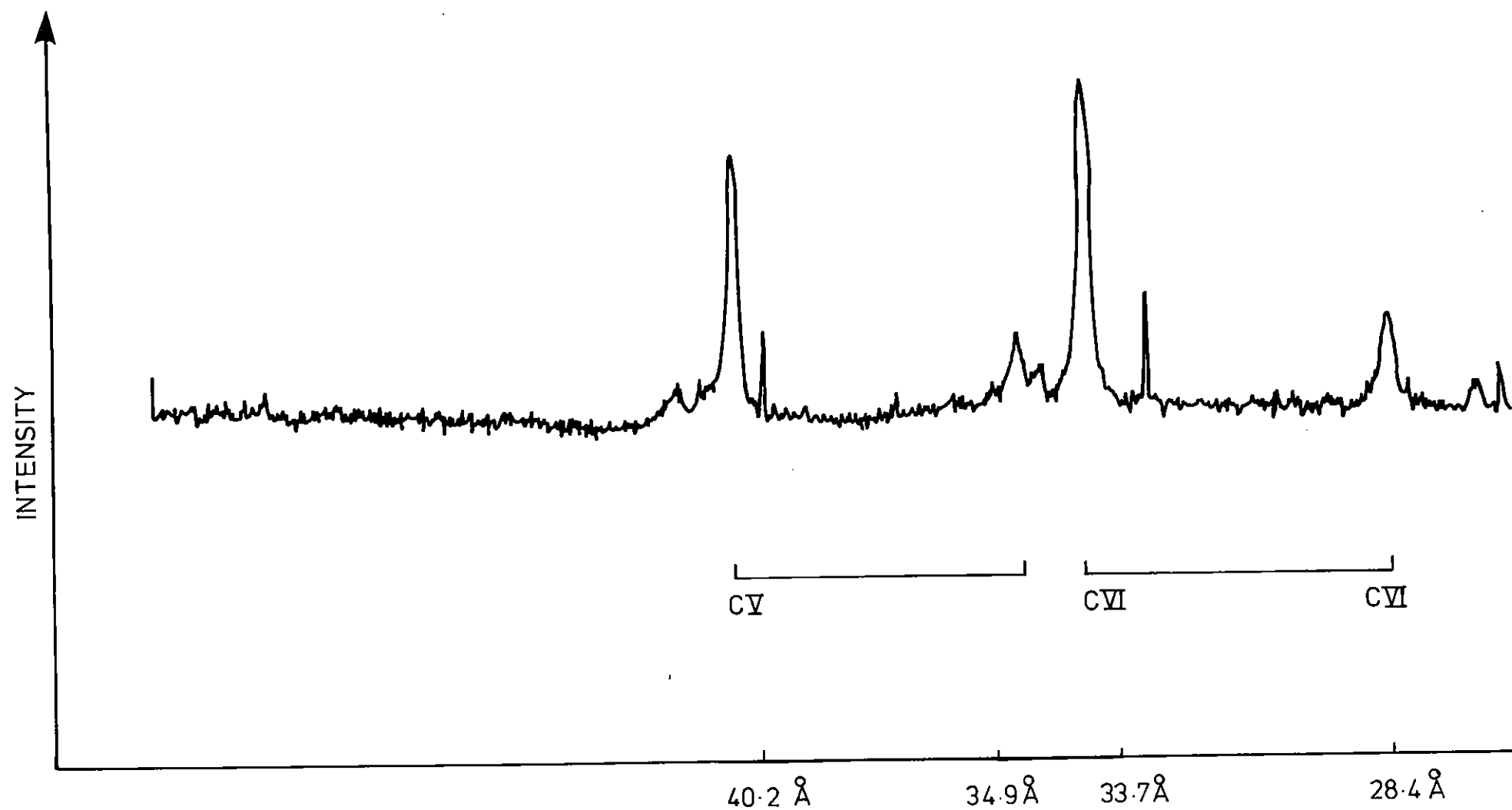


Fig. A4.33 Carbon line spectra recorded on rocket grazing incidence grating spectrograph.

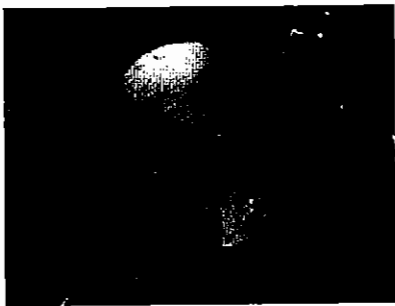


Fig. A4.34 Image plane of zone plate.

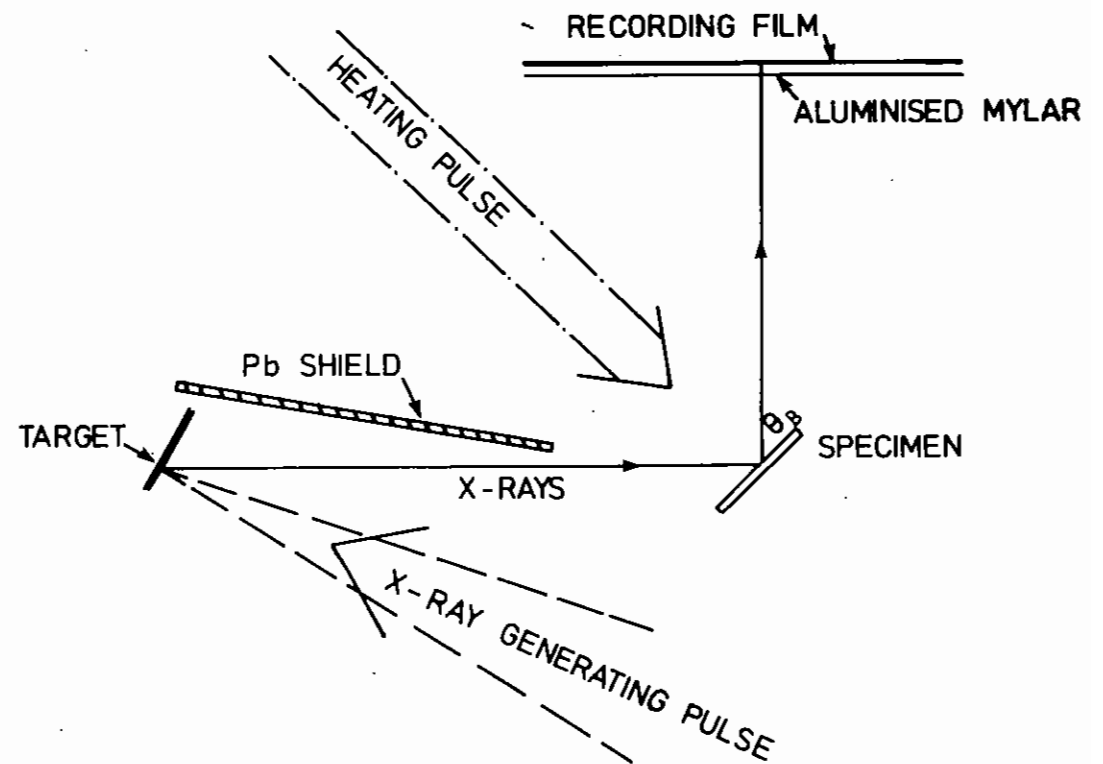
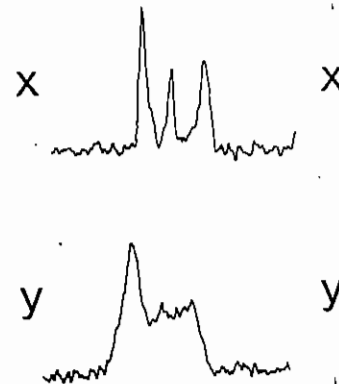
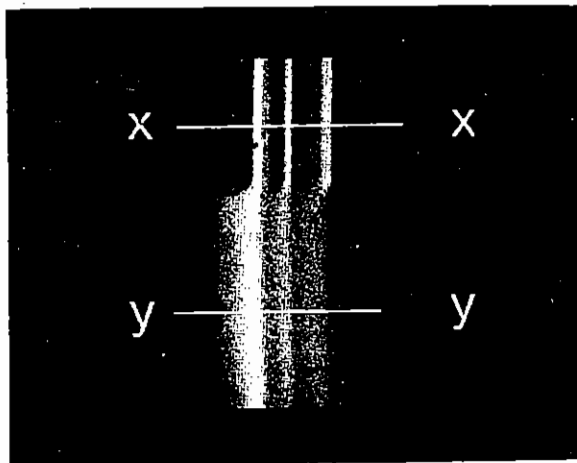


Fig. A4.35 Experimental X-ray diffraction apparatus.

TIME RESOLVED LASER ANNEALING OF SILICON (111)

TIME RESOLVED LASER ANNEALING OF SILICON (111)

SHOT 19 10/1/84



CHLORINE He α
LINE FROM PVC TARGET

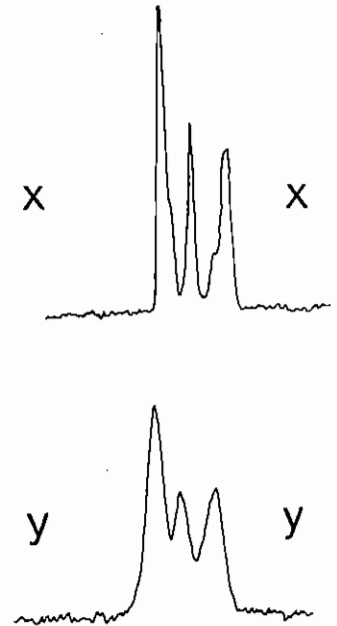
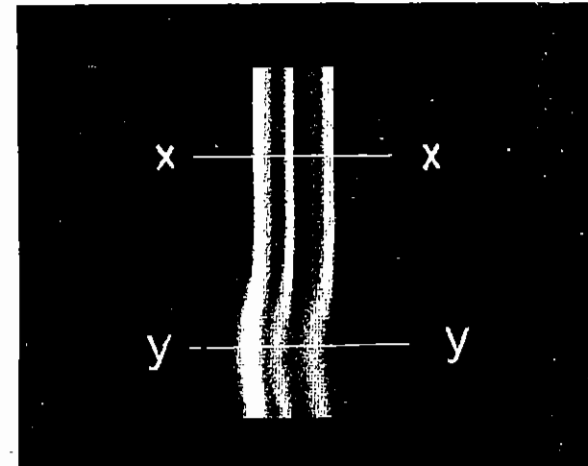
LASER ENERGY IN ANNEALING BEAM $\sim 1 \text{ J/CM}^2$

LASER ENERGY FOR XRAY DIFFRACTION PROBE 15 J

TIME DELAY BETWEEN ANNEALING PULSE AND DIFFRACTION PROBE 80 NS

Fig. A4.36 Spectrum from Silicon (111) recorded at 80ns after heating pulse.

SHOT 15 13/1/84



CHLORINE He α
LINE FROM PVC TARGET

LASER ENERGY IN ANNEALING BEAM 1-2 J/CM²

LASER ENERGY FOR XRAY DIFFRACTION PROBE 22 J

TIME DELAY BETWEEN ANNEALING PULSE AND DIFFRACTION PROBE 182 NS

Fig. A4.37 Spectrum from Silicon (111) recorded at 182ns after heating pulse.

temperature-depth relation in the outermost few μm .

(c) The changes to the diffraction patterns are as a consequence of (b) very marked compared with the data of Larson et al.

The principal disadvantage is that we are working at a wavelength where the scattering factor and absorption length are less well known. This creates a problem in that our sampling depth is not very well defined. Since we are diffracting off a high quality single crystal we also have the usual problems with extinction effects. If we use the absorption coefficient to estimate our sampling depth we have a figure of $\sim 3.5 \mu\text{m}$, whereas if we use the extinction coefficient we arrive at a figure of $\sim 0.34 \mu\text{m}$. Actually, a similar problem should also occur with the analysis of the synchrotron data. In order to further probe the near surface region we adopted a low angle of incidence geometry as shown in Figure (A4.38). We also used a silicon (100) surface and oriented the crystal so that we diffracted off (111) planes, this time using $\lambda = 5.68 \text{ \AA}$. This greatly reduces our sampling depth to between $0.78 \mu\text{m}$ and $0.34 \mu\text{m}$ depending on whether we use absorption or extinction coefficients. The results confirmed this since we saw a broadening and shift of our diffracted beams at much shorter times than in the other experiment.

The analysis of the results is proceeding as follows. We set up a model of our X-ray beam interaction with the silicon and assume an attenuation law in the material, either excluding or including extinction effects. A notional temperature distribution is then assumed, and the diffraction angles and intensities are calculated and compared with experiment. Adjustments of the temperature distributions are made until agreement is obtained. This phase of the work is incomplete, but the results appear to be in broad agreement with those of Larson et al (A4.35) especially with regard to the delay times for producing surface melting and the cooling rates.

These preliminary experiments have very successfully demonstrated the potential use of the pulsed laser plasma X-ray source at RAL for probing the changes in crystallography occurring in the near surface region of material on a time resolved basis of $\sim \text{nsec}$. It should be possible to extend this to

less than 100 ps without too much difficulty. This opens the way to study a wide range of important problems encountered in laser processing of semiconductors, eg amorphisation or recrystallisation, interface compound formation (silicides) redistribution of implanted impurities etc. There are also many other very interesting solid state transformations (eg martensitic transformations) which will be amenable to study with this technique.

A4.7 Energy Loss of Charged Particles in Dense Plasmas

A P Fewes, D L Henshaw (Bristol), S P Knight, D A Pepler and W T Toner (RAL)

In the early part of the year a short run was devoted to improving techniques for producing exploding pusher implosions. No single cause for the disappointing results of some earlier runs was identified although prepulse was found to be somewhat larger than had been thought and the targets and pulse length then available were not well matched to the energy. Following improvements in all three areas, neutron yields in excess of 10^7 were again obtained from four-beam implosions in the spring of 1984. A four week run was then devoted to systematic data collection on the energy loss of charged particles in dense plasmas (A4.37). CR39 exposures to α particles traversing mylar, aluminium, copper and gold foils were made for several different heating conditions and protons and tritons were recorded in a more limited set of conditions. It will be some months before a quantitative assessment of the data is available. Automatic track measurement techniques are discussed in A6.

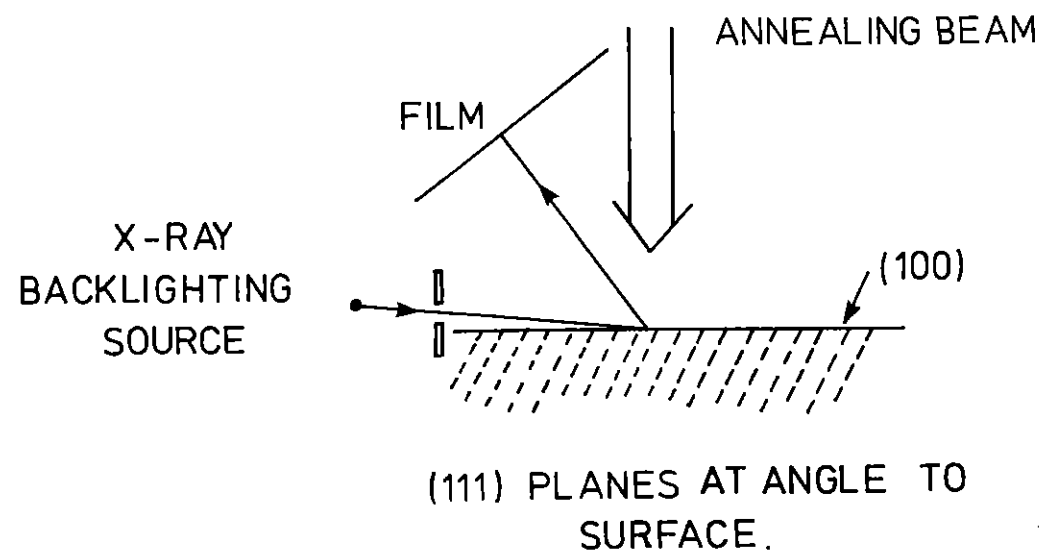


Fig. A4.38 Low angle of incidence geometry for diffraction from (111) planes.

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A5 THEORY AND COMPUTATION

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A5 THEORY AND COMPUTATION

A5.1 Introduction

R A Cairns (St Andrews)

Theoretical work on plasma physics continues that of previous years in two main areas, studies of interaction and absorption processes and of the kinetic theory of transport processes. In the first a particle code has been used to look at some of the effects taking place near the quarter critical surface and a number of studies have considered the interaction between different processes occurring in the plasma. Kinetic theories of heat flow, which have been a very active field of research recently, have been extended to take account of ablation, while work on heat flow instabilities has been extended to a kinetic rather than a fluid model.

In hydrodynamics the development of computer codes continues together with the use of these codes and other means to study instabilities of a compressed target. A substantial effort has also gone into theoretical studies of ion beam irradiated targets. The theory of dense plasmas is a subject which is attracting more attention as the experimental facility becomes increasingly capable of investigating this regime, and a number of contributions investigate radiation transport and atomic physics effects in such plasmas. Finally there is a contribution to laser physics, with a study of the behaviour of a backward Raman amplifier.

A5.2 Plasma Physics

A5.2.1 Effects of Linear Mode Conversion on Raman Scatter

G A Gardner, H C Barr, T J M Boyd and R Rankin (Bangor)

Some attention has been paid to the role which self-generated magnetic fields may play in parametric processes (A5.1 - A5.3) and in particular Boyd (A5.2) and Estabrook and Kruer (A5.3) have suggested that the megagauss fields characteristic of certain laser target experiments may affect

back-scattered Raman light significantly. These fields may give rise not only to measurable frequency shifts but can alter significantly the levels of Raman backscatter expected from the plasma. In this report we consider simulations of Raman backscatter from an inhomogeneous plasma in the presence of strong self-generated magnetic fields. The emission levels show a dependence on magnetic field which is consistent with the scattered laser light undergoing linear mode conversion.

In order to investigate the frequency spectrum and reflectivities of Raman scattered light we have performed a series of computer simulations using a 1.5 dimensional fully relativistic electromagnetic particle code. In all the simulations the incident laser light, with an intensity equivalent to an electron oscillatory velocity of $v_0 = 0.1c$ ($1.2 \times 10^{16} \text{ Wcm}^{-2}$ for $1.06 \mu\text{m}$ light), propagated through a plasma with density profile rising from zero to $0.35N_c$ in a distance of either $6.7\lambda_0$ or $19\lambda_0$ where λ_0 is the laser wavelength. The initial electron temperature was 1.3 KeV . Ion motion was included with $m_i/m_e = 1000$ and initially $T_i/T_e = 2$. This choice of temperature ratio did not wholly eliminate Brillouin scatter but kept it at a low level. Constant external magnetic fields were applied parallel to the laser light magnetic field with electron cyclotron frequencies Ω in the range 0 to $0.1\omega_0$ (ω_0 is the laser frequency), i.e., up to 10 MG for $1.06 \mu\text{m}$ laser light. Thus, the laser light propagated as an extraordinary wave and decayed into an upper hybrid wave and a backscattered extraordinary wave.

Fig.A5.1 shows frequency spectra obtained from the simulations for $\Omega = 0.075 \omega_0$ and $\Omega = 0.0375 \omega_0$. The frequency of back scattered light is plotted against Ω in Fig.A5.2 from which it is clear that the characteristic thermal red shift (A5.4) of the Raman light wave can be changed to a blue shift for sufficiently large magnetic fields.

An estimate of this magnetic shift may be made as follows. Since $\Omega \ll \omega_0$, to first order in Ω , we may treat both the laser light and the upper hybrid wave as unmagnetized whereas the scattered extraordinary wave emerging from its reflection point near $N_c/4$ (where the instability is absolute and the threshold a minimum) has a frequency ω_s given by $\omega_s^2 = \omega_p^2 + \omega \Omega$. Phase matching yields simply that

A5.2

$$\omega_s = \frac{\omega_o}{2} - \frac{9}{8} \frac{V_e^2}{c^2} + \frac{\Omega}{4} \quad (1)$$

where V_e is the electron thermal velocity. Thus, the thermal red shift is countered by a blue shift, $\Omega/4$, due to the magnetic field. In terms of wavelength, this magnetic field-induced effect may be described by $\Delta\lambda_s(\text{\AA}) = 93.3\lambda_o^2(\mu\text{m}) B(\text{MG})$. The simulations show broad agreement with this with the exception of the red wing of the double peaked structure appearing in Fig. A5.1(b) and in Fig. A5.2 when $\Omega = 0.0375\omega_o$. We shall return to this below.

The reflectivities for the two different scale lengths simulated are shown in Fig.A5.3 as a function of magnetic field. These are time integrated over the whole run although in practice the emission does not switch on until well into the run and then persists only long enough for the density profile to steepen sufficiently to saturate the instability. Pulsing of the emission is observed for low magnetic fields which is consistent with profile steepening followed by relaxation allowing a further burst of emission. This disappears at the larger magnetic field values. Fig.A5.3 may be understood in terms of linear mode conversion of the scattered extraordinary wave, at frequency ω_s , into an upper hybrid wave.

The linear mode conversion of a light wave to a plasma wave is well understood as resonance absorption. Efficient conversion requires the resonance of the plasma wave to be in close proximity to the reflection point of the light wave and in addition, a finite electric (driver) field in the direction of the density gradient. This, in turn, requires waves propagating obliquely to the density gradient - a two dimensional problem. Alternatively, a D.C. magnetic field allows coupling in one dimension (A5.5). The critical parameter is $\tau = (\omega L/c)^{1/3} (\Omega/\omega)^{1/2}$ where ω is the frequency of the wave involved and L is the local density scale length (or distance from the vacuum - plasma boundary - to $N_c/4$ in the case of the linear density profile assumed here). Maximum conversion and hence absorption occurs when $\tau \approx 0.8$. We note also that the reverse process, namely conversion from the upper hybrid wave, at frequency $\omega_o - \omega_s$, to an extraordinary wave, exhibits the same dependence on τ . Consequently, both

processes are correspondingly weak or strong and provide additional damping of the daughter waves of the instability, thereby increasing the threshold.

Fig.A5.3 shows for each of the scale lengths considered that the emission is substantially reduced as the magnetic field increases and is a minimum for some critical field strength. This behaviour reflects the increasing value of τ and hence the increased absorption of the scattered wave by mode conversion. The minimum reflectivities occur where $\tau \approx 0.8$ and mode conversion is a maximum. In fact, for the case with the initially steep profile (and hence nearer to instability threshold), the emission is effectively switched off for a range of magnetic fields before recovering to about its unmagnetized value. For the gentler density gradient, the instability was of course much stronger showing much higher emission levels and hence the behaviour was easier to diagnose. These simulations showed severe profile steepening for low magnetic fields and again for high magnetic fields. In the intermediate regime where mode conversion was optimum, the steepening was rather moderate suggesting that a combination of profile steepening and mode conversion now saturates the instability. A third regime occurs for this case. The reflectivity has not recovered to its initial level but levels off at about 3.5%. However, for these magnetic fields, strong steepening is observed thereby reducing the value of τ and increasing the absorption by mode conversion.

Finally, we remark that the inverse process of conversion from upper hybrid wave to extraordinary wave at a frequency $\omega_o - \omega_s$ is optimum also around $\tau \approx 0.8$. This would appear to explain the red wing observed in Fig.A5.1(b) and in Fig.A5.2 when $\Omega = 0.0375 \omega_o$. For other values, either τ was never optimum or not optimum sufficiently long for this feature to be observed.

In conclusion, the simulations demonstrate that self-generated magnetic fields result not only in observable frequency shifts (100 Å per megagauss for 1μm light) but also in reduced levels of Raman backscatter. We attribute these reduced levels to linear mode conversion of the daughter waves of the Raman instability.

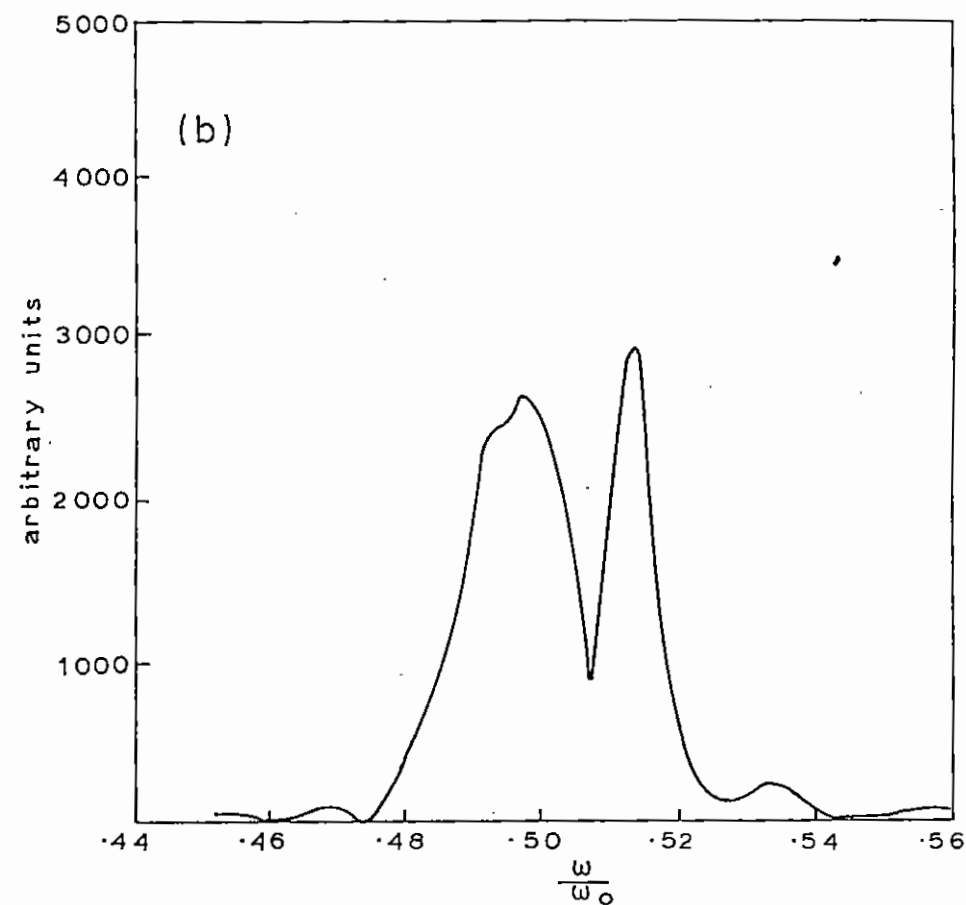
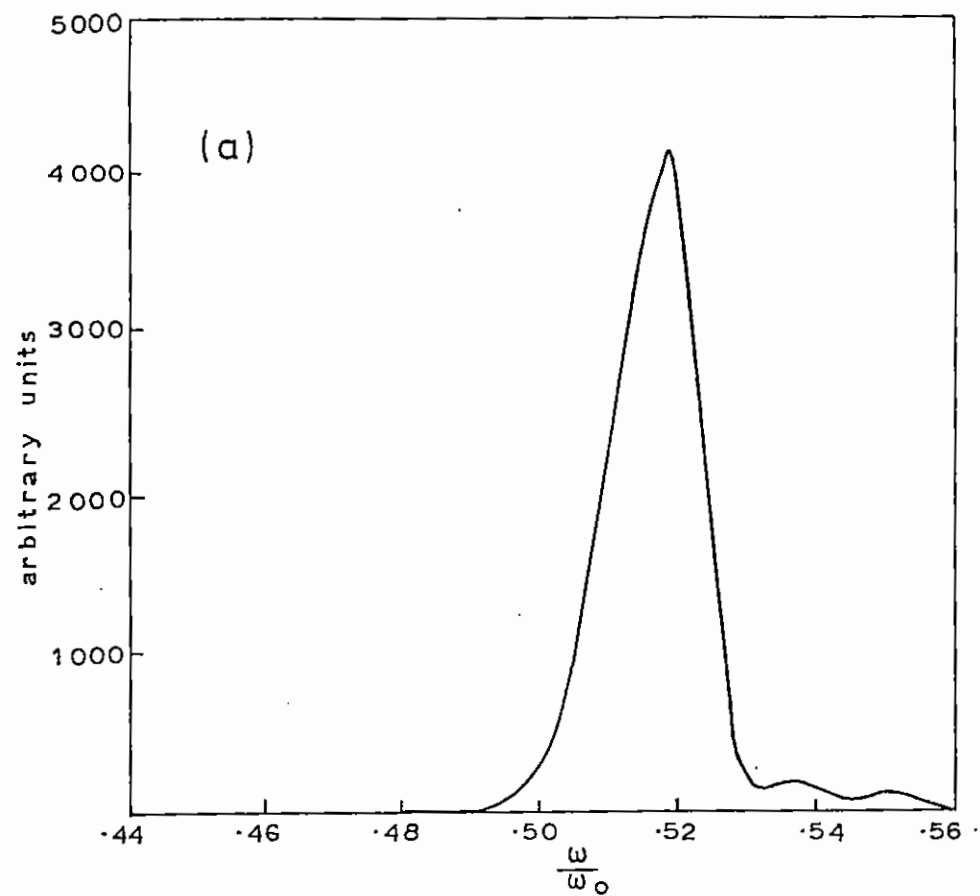


Fig. A5.1 Frequency spectrum of the backscattered light measured in simulations; $v_0/c = 0.1$, $V_e/c = 0.1$ with $\Omega = 0.075\omega_0$ for Fig A5.1(a) and $\Omega = 0.0375\omega_0$ for Fig A5.1(b).

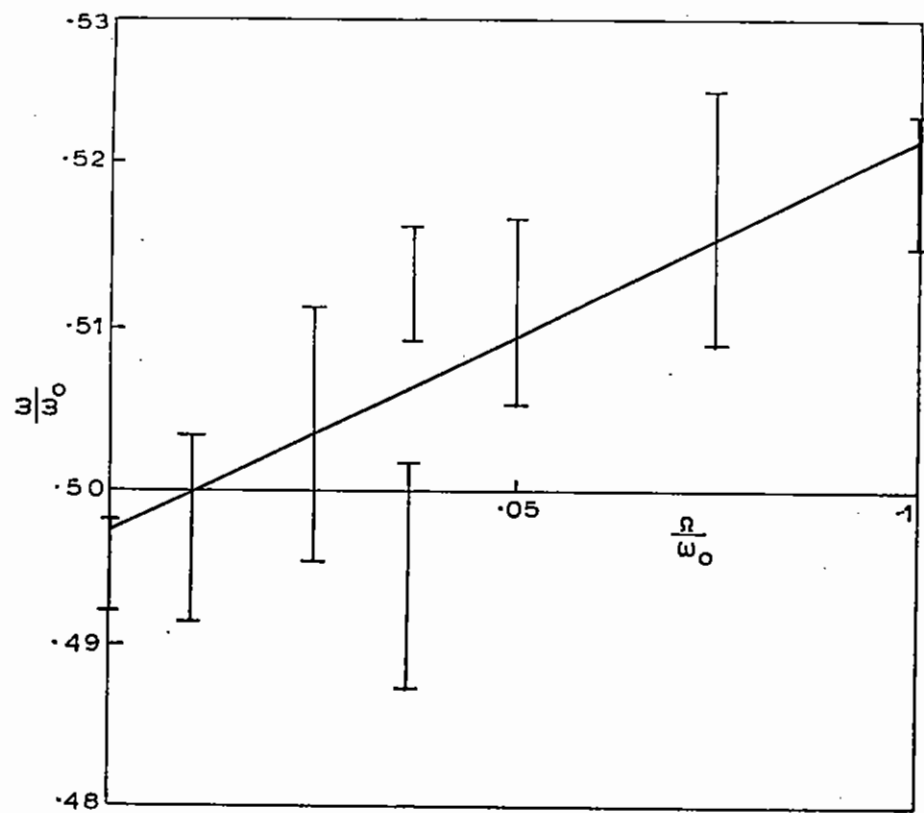


Fig. A5.2 Frequency of the backscattered light as a function of Ω . The solid line represents the theoretical values from Eq. 1 and the results of the simulations are represented by I.

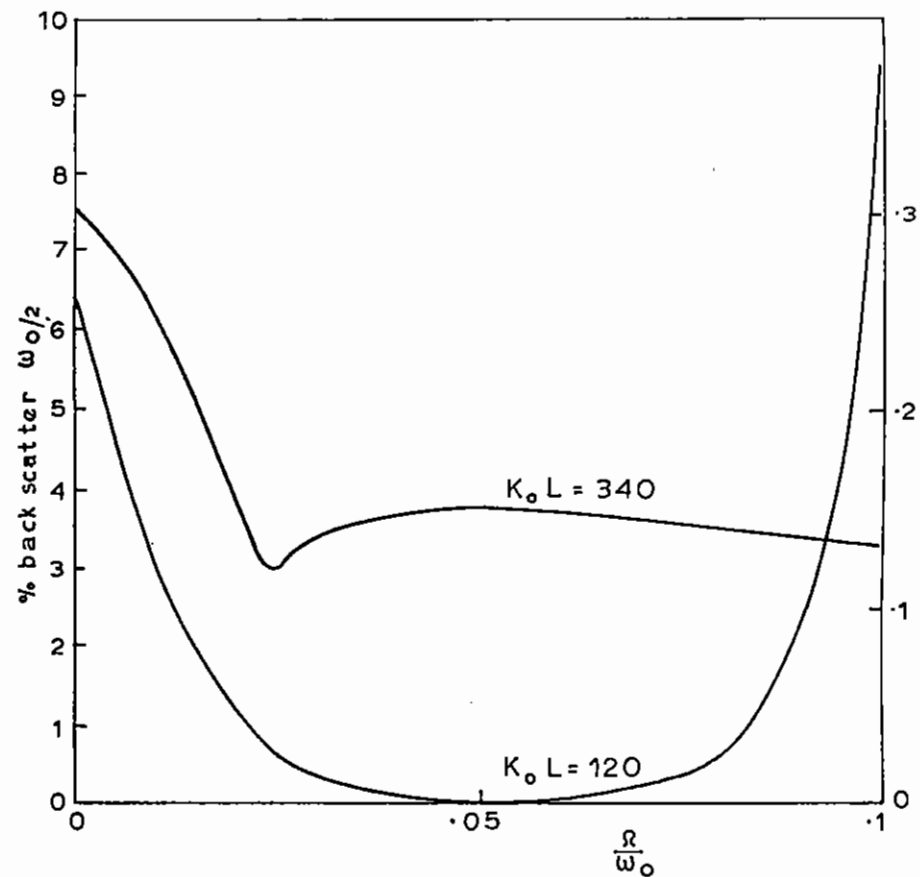


Fig. A5.3 Reflectivity plotted vs. Ω , for two scale lengths. The left hand vertical scale corresponds to $k_0 L = 340$; that on the right hand side to $k_0 L = 120$.

A5.2.2 Linear Mode Conversion in a Magnetised plasma

G A Gardner, R Rankin, H C Barr and T J M Boyd (Bangor)

Resonance absorption, the linear mode conversion of a light wave to a plasma wave, has been extensively studied as an absorption mechanism in laser target experiments. The inverse process, whereby a plasma wave may convert to a light wave, has received less attention in this context. Both are particularly relevant to processes occurring at the quarter-critical density of laser produced plasmas by providing additional damping of the daughter waves of parametric instabilities thus, for example, reducing the level of Raman scatter and giving rise to additional emission as a result of radiation from plasma waves. Efficient conversion requires the close proximity of the plasma wave resonance and the electromagnetic wave reflection point, as well as a finite (driver) electric field in the density gradient direction. This in turn requires either waves propagating obliquely to the density gradient (a two dimensional problem), or alternatively a DC magnetic field which allows coupling in one dimension (A5.4). It is this problem we examine here, namely the linear mode conversion of an upper hybrid wave to an extraordinary wave and vice-versa.

The equations describing the propagation of both extraordinary and upper hybrid waves in an inhomogeneous plasma with a uniform applied magnetic field $\underline{B}_c = B_c \hat{z}$ are

$$c^2 E''_y + (\omega^2 - \omega_p^2) E_y = i \omega \omega_c E_x$$

$$3V_e^2 E''_x - 3V_e^2 L^{-1} E'_x + (\omega^2 - \omega_p^2 - \omega_c^2) E_x = -i \omega^{-1} \omega_c \omega_p^2 E_y$$

where L is the density gradient scale length, the profile being assumed linear. These equations may be routinely solved with boundary conditions chosen to represent an incident upper hybrid wave or an incident extraordinary wave. The conversion efficiency for both cases may then be determined. Both show the expected (A5.4, A5.6) dependence on $\tau = (\omega L/c)^{1/3} (\omega_c/\omega_0)^{1/2}$ with optimum conversion efficiency of around 60% when $\tau \approx 0.8$.

The particle simulations were performed using a 1½D relativistic electromagnetic code.

An upper hybrid wave was excited at its resonance by applying a weak locally driven electric field along the density gradient. No incident laser beam was present, the ions were fixed, and an electron temperature corresponding to $V_e/c = 0.05$ was used. The variation of conversion efficiency with magnetic field strength and scale length was investigated and the results are summarised in Fig.A5.4. A maximum conversion of about 30% at the predicted value of $\tau (\approx 0.8)$ was observed. The conversion efficiency was determined as a time integrated value over the whole time of a run; time integration over only that portion during which significant emission was observed gave values $\approx 50\%$ in reasonable agreement with the predicted theoretical values.

A typical run shows the growth of the upper hybrid wave, its saturation by wave breaking at about 10 times the driver field, and the conversion to extraordinary waves which propagate out of the simulation region. Some second harmonics of the driver frequency were observed in the light spectra emerging from the plasma.

We have demonstrated that linear mode conversion of plasma waves to light waves can occur in a magnetised plasma. The process can be efficient (efficiencies of up to 50% can be attained), and scales with $\tau = (kL)^{1/3} (\omega_c/\omega)^{1/2}$ in the expected manner. In Raman scatter near quarter-critical density, linear mode conversion of each of the daughter waves provides additional damping resulting in reduced levels of emission.

A5.2.3 Kinetic model for the collisional Weibel instability in laser-produced plasmas.

E M Epperlein (Imperial College)

Recently there have been many experimental reports on the occurrence of filamentary structures (or plasma jets) and strong magnetic fields in the

ablation of laser driven targets.(A5.7,A5.8) These structures, which have transverse scale-lengths of about 10 μ m, occur mainly in medium to high Z targets (e.g. glass or gold), and have been inferred to originate at densities above critical for 1 μ m laser light.(A5.7) There, they may have serious effects on the target implosion symmetry, by inhibiting lateral transport and inducing pressure perturbations in the ablation front.

Several mechanisms have been proposed to explain the appearance of these jets, such as the radiation cooling instability,(A5.9) the thermal instability(A5.10) and the collisional Weibel instability.(A5.8) Here, the last mechanism is investigated, as it most naturally leads to the formation of current filaments and magnetic field self-generation predicted from the experiments.(A5.8)

The model for the collisional Weibel instability developed by Mochizuki et al.(A5.8) consists of artificially splitting the electrons into hot and cold fluids, in a homogeneous background of cold and fixed ions. Using planar geometry, with the target parallel to the xy plane, the hot electrons are assumed to drift in the positive z direction with a speed v_h and the cold electrons in the opposite direction with a speed v_c . In order to preserve charge neutrality, it is required that $n_h v_h = n_c v_c$ and $Z n_i = n_h + n_c$, where n is number density and Z the ionic charge (subscripts c,h and i refer to cold electrons, hot electrons and ions, respectively). This configuration leads to a heat flow along z given by: $q = \frac{1}{2} m(n_h v_h^3 - n_c v_c^3) + \frac{5}{2} (n_h T_h v_h - n_c T_c v_c)$, where m is the electron mass and T the temperature (in energy units). This configuration also gives rise to a pressure anisotropy which acts as a source to drive the unstable Weibel mode.(A5.11)

The instability mechanism is briefly described as follows (A5.8). An induced magnetic field in the y direction, with a modulation along x, pushes the electrons towards its neutral points, via the Lorentz force. There, a net current is created (via Ampere's law), which is directed so as to amplify the initial magnetic field perturbation, thus promoting an instability. Its growth rate will be positive unless damped out by Joule dissipation, at short perturbation wavelengths.

A kinetic model for this instability will now be developed. The equilibrium state is assumed equivalent to the one just described, and is represented by a two temperature electron distribution (with $Z n_i = n_c \gg n_h$) defined as follows.

$$f = \left\{ (n_c/v_{Tc}^3) \exp(-(v_z + v_c)^2/v_{Tc}^2 - v_{\perp}^2/v_{Tc}^2) + (n_h/v_{Th}^3) \exp(-(v_z - v_h)^2/v_{Th}^2 - v_{\perp}^2/v_{Th}^2) \right\} / \pi^{3/2}$$

where $v_T = (2T/m)^{1/2}$ is the electron thermal speed and $v_{\perp}^2 = v_x^2 + v_y^2$.

The time evolution of the plasma is governed by the electron Fokker-Planck equation (introducing the notation $\partial_w = \partial/\partial w$, for a coordinate w),

$$\partial_t f + \underline{v} \cdot \nabla f - (e/m) (\underline{E} + \underline{v} \times \underline{B}/c) \cdot \nabla_v f = C,$$

and Maxwell's equations,

$$\partial_t \underline{B} = -c \nabla \times \underline{E}, \quad \nabla \times \underline{B} = 4\pi \underline{j}/c,$$

where e is the electronic charge, c the speed of light, $\underline{j} = \int d\underline{v} \underline{v} f$ the electric current, E and B the electric and magnetic field intensities, respectively, and C the collision operator (A5.12). The displacement current has been dropped from Ampere's law, since we may neglect relativistic effects.

Before linearizing the above equations, it is found convenient to expand the distribution f in cartesian tensors:

$$f = f_0 + \underline{v} \cdot \underline{f}_1/v + \frac{\underline{v} \underline{v} \cdot \underline{f}_2}{v^2} + \frac{\underline{v} \underline{v} \underline{v} \cdot \underline{f}_3}{v^3} + \dots$$

This expansion is expected to converge provided the collisional term on the R.H.S. of the Fokker-Planck equation dominates over all other terms.(A5.13) Since we are only interested in collisional phenomena we keep terms up to

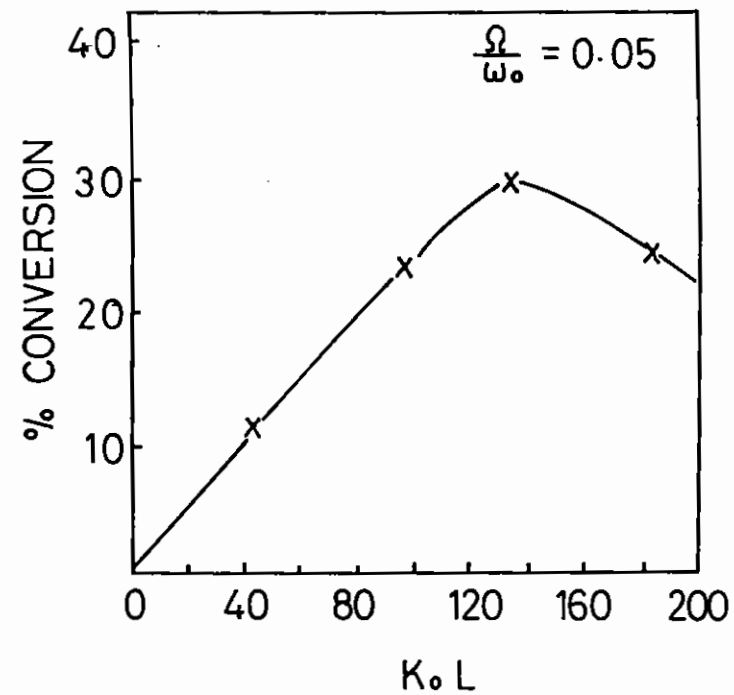
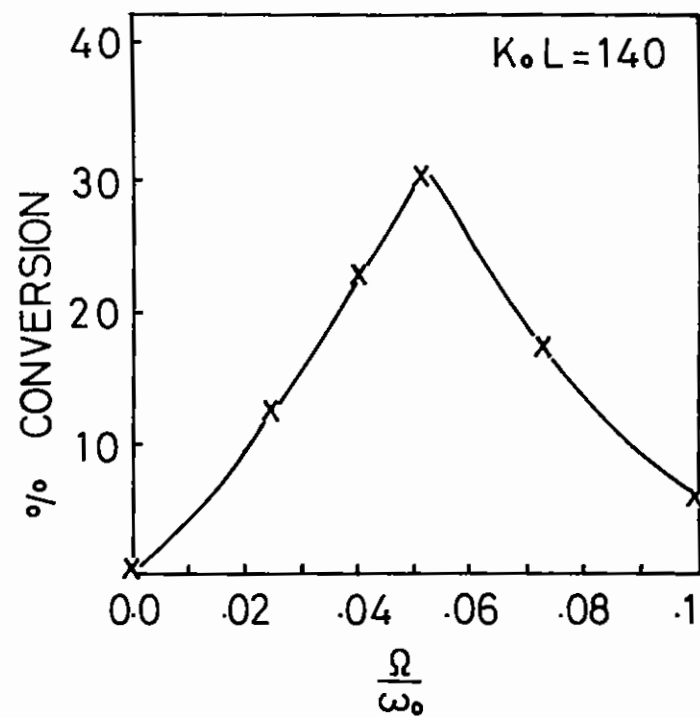


Fig. A5.4 The variation of conversion efficiency with magnetic field strength and scale length.

and including $\underline{f}_3$, where $f_0 \gg \underline{f}_1 \gg \dots$ is assumed. Substituting this truncated expansion into the Fokker-Planck equation, and making use of its orthogonal properties, one arrives at four coupled equations in f_0 , $\underline{f}_1$, $\underline{f}_2$ and $\underline{f}_3$ (A5.13). Linearizing these and Maxwell's equations, and taking a perturbation of the form $\exp(\gamma t + i k x)$, we obtain (with the prime denoting perturbed quantities):

$$\gamma f'_0 + \frac{1}{3} k v f'_x - \frac{1}{3} \epsilon'_z \left[(1/v^2) \partial_v (v^2 f'_z) \right] = C'_{ee0}, \quad (1)$$

$$i k v f'_0 - \epsilon'_x \partial_v f'_0 - \omega' f'_z = -v_{ei} f'_x + C'_{eex}, \quad (2)$$

$$\frac{2}{5} i k v f'_{xz} - \epsilon'_z \partial_v f'_0 = -v_{ei} f'_z + C'_{eez} \quad (3)$$

$$\begin{aligned} \frac{1}{2} i k v f'_z + \frac{3}{7} i k v f'_{xxz} - \frac{1}{2} \epsilon'_x (v \partial_v (f'_z/v)) - \frac{1}{3} \omega' f'_{zz} \\ = -3 v_{ei} f'_{xz} + C'_{eexz}, \end{aligned} \quad (4)$$

$$\frac{8}{15} i k v f'_{xz} = -6 v_{ei} f'_{xxz} + C'_{eexxz}, \quad (5)$$

$$\gamma \omega' = i k \epsilon'_z, \quad (6)$$

$$i k \omega' = -(\omega_p/c)^2 (4\pi/3 n_c) \int_0^\infty dv v^3 f'_z, \quad (7)$$

where $\underline{\epsilon}' = (e/m)(\underline{E}'_x, 0, \underline{E}'_z)$, $\underline{\omega}' = (e/mc)(0, B', 0)$, $\omega_p = (4\pi n_c e^2/m)^{1/2}$ is the plasma frequency, $v_{ei} = (4\pi n_i (Ze^2/m)^2 \ln \Lambda_{ei})/v^3$ the electron-ion collision frequency and $\ln \Lambda_{ei}$ the Coulomb logarithm.

In deriving the above equations it has been assumed that $\gamma, kv \ll v_{ei}$, for a thermal electron, in accordance with the collisionality condition. The effect of electron self-collisions in the $\underline{f}_1$, $\underline{f}_2$ and $\underline{f}_3$ equations may be approximately taken into account by rewriting $\underline{v}_{ei}$ as $\underline{v}_{ei} + v_{ee} = v$, where $v_{ee} = (4\pi n_c (e^2/m)^2 \ln \Lambda_{ee})/v^3$.

Without further approximations, equations (3)-(7) give rise to the following dispersion relation (written in dimensionless form):

$$\begin{aligned} \Gamma = \left\{ \int_0^\infty \left(\frac{a K^2 v^{10}}{1+(a+b)K^2 v^8} \right) (3F_{zz} + v \partial_v (F_z/v)) dv - 3K^2 L_s^2 \right\} \\ \left\{ - \int_0^\infty \left(\frac{1+bK v}{1+(a+b)K^2 v^8} \right) v^6 \partial_v F_0 dv \right\}^{-1}, \end{aligned} \quad (8)$$

where $a = 1/15$, $b = 4/315$, $\Gamma = \gamma/v_{Tc}$, $K = k\lambda_{Tc}$ ($=kv_{Tc}/v_{Tc}$), $L_s = (c/\omega_p)/\lambda_{Tc}$, $v = v/v_{Tc}$, $F = (4\pi/n_c) v^3 f$, $\phi = \epsilon'_x/\omega' v_{Tc}$ and $v_{Tc} = v(v_{Tc})$. The term ϕ , which relates ϵ'_x to ω' through the quasi-neutrality condition

$$\int_0^\infty dv v^2 f'_0 = \int_0^\infty dv v^3 f'_x = 0,$$

should strictly be evaluated by solving equations (1), (2) and (6).

Instead, we derive a lower and upper bound for ϕ , i.e. ϕ_l and ϕ_u , by solving equation (2) in the limit of $ikf'_0 \gg v f'_x$ and $ikf'_0 \ll v f'_x$, respectively, and investigate their overall effect on the results.

Equation (8) is now solved numerically for the following test problem: $n_h/n_c = 0.005$, $T_h/T_c = 3$, $v_c = 0.01$ and $v_h = 2$. These parameters are consistent with typical Fokker-Planck simulations of transport in plasmas irradiated by $1\mu m$ laser light (A5.14, A5.15). They represent a plasma with a heat flow which is about 13% of the 'free-streaming' limit $q_f = n_c m (T_c/m)^{3/2}$

and a cold electron drift at the isothermal sound speed $c_s = (2T_c/m)^{1/2}$ (for $Z=10$).

In Fig. A5.5 we plot Γ as a function of Λ ($=2\pi/K$) for the case of no short-wavelength damping, i.e. for $L_s = 0$. The solid lines indicate the solution of equation (8), for $\phi = 0$, where the one labelled 'a' refers to case in which $\underline{f}_3$ is included in the analysis, and the one labelled 'b' refers to case in which it is not (obtained by simply setting $b=0$ in equation 8). As observed, the contribution due to $\underline{f}_3$ decreases for large Λ , which is expected since the magnitude of the spatial advection term in the

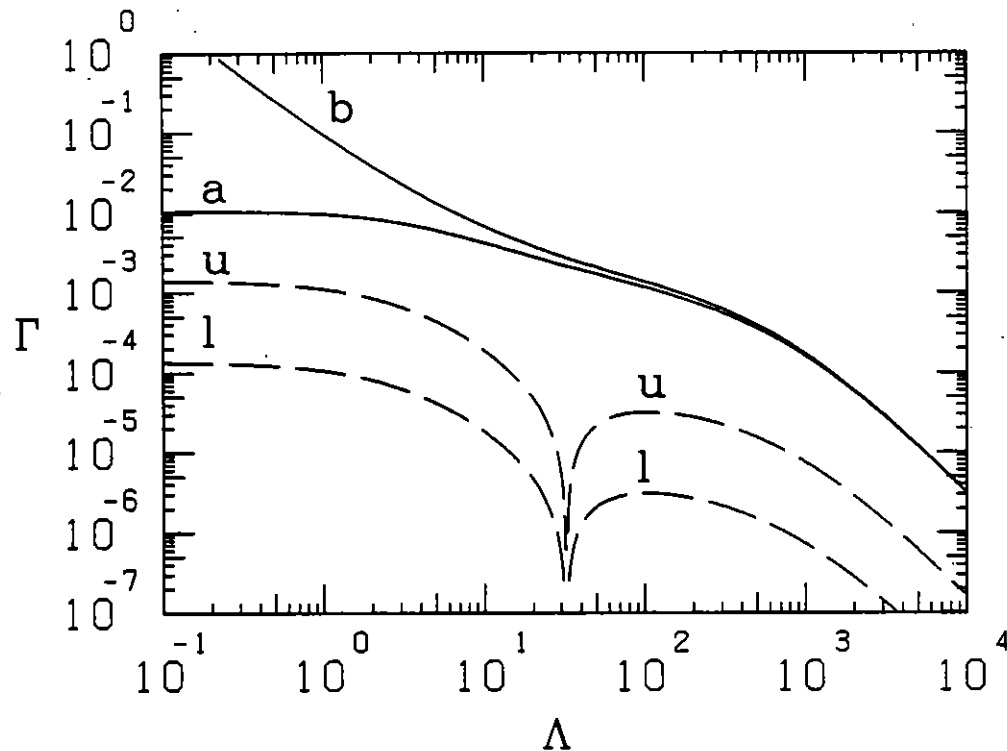


Fig. A5.5

Plot of growth rate Γ vs. perturbation wavelength Λ , for $L_s = 0$, $n_h/n_c = 0.005$, $T_h/T_c = 3$, $v = 0.01$ and $V_e = 2$ ($q/q_e = 0.13$). Solid lines are the results for $\phi = 0$, with: (a) f'_1 contribution and (b) no f'_1 contribution. Dashed lines show the magnitude of the contributions due to: (l) ϕ_l and (u) ϕ_u .

Fokker-Planck equation becomes small in comparison with the collision term. In the opposite limit ($\Lambda \ll 1$) the expansion in f' starts to break down, indicating that collisionless phenomena become important. The dashed curves represent the magnitude of the contributions due to ϕ_l (labelled 'l') and ϕ_u (labelled 'u'), which, as observed, are negligible.

In Fig. A5.6 we see the effect on the growth rate of introducing various amounts of damping, by letting $L_s = 0.05$, 0.5 and 5 . For comparison, the results obtained from the collisionless Weibel theory (A5.16) (dotted lines) and the 2-fluid collisional theory of Mochizuki et al. (A5.8) (dashed lines) are also presented. Clearly, the collisionless theory inadequately predicts the growth rate when $\Lambda > 1$. As for the 2-fluid theory, it underestimates the growth rate at very short wavelengths and has an incorrect K dependence at long wavelengths.

An important feature of the present theory is the presence of a minimum value of L_s above which no unstable Weibel mode can occur. From equation (8) this critical L_s is given by:

$$L_s^c = 5.48 \times 10^{-2} (n_c/10^{21} \text{ cm}^{-3})^{1/2} (Z/10) (\ln A_{ei}/5) (1 \text{ keV}/T_c)^2 \\ = (a \int_0^\infty dv v^{12} F_{zz})^{1/2},$$

Hence, for a cold and dense high- Z plasma, Joule dissipation will damp out any B-field generation by the Weibel instability. For the equilibrium configuration under study, $L_s^c = 5$ (as may be verified from the plot in Fig. A5.6), which implies that for $Z=10$ and $n_c = 10^{23} \text{ cm}^{-3}$ $T_c > 0.3 \text{ keV}$ is a necessary condition for instability.

In an attempt to correlate these theoretical predictions with the experimental observations (A5.7, A5.8), we plot the maximum growth rate in Fig. A5.7(a) and corresponding perturbation wavelength in Fig. A5.7(b) as a function of Z , for several values of T_c and n_c . From Fig. A5.7(a) we note that, in general, the growth rate increases with Z , in agreement with experiment (A5.7). The optimum wavelength, shown in Fig. 3(b), is relatively insensitive to variations in n_c , T_c and Z , and compares well with the observed spatial scale-lengths of the plasma jets. It must be stressed,

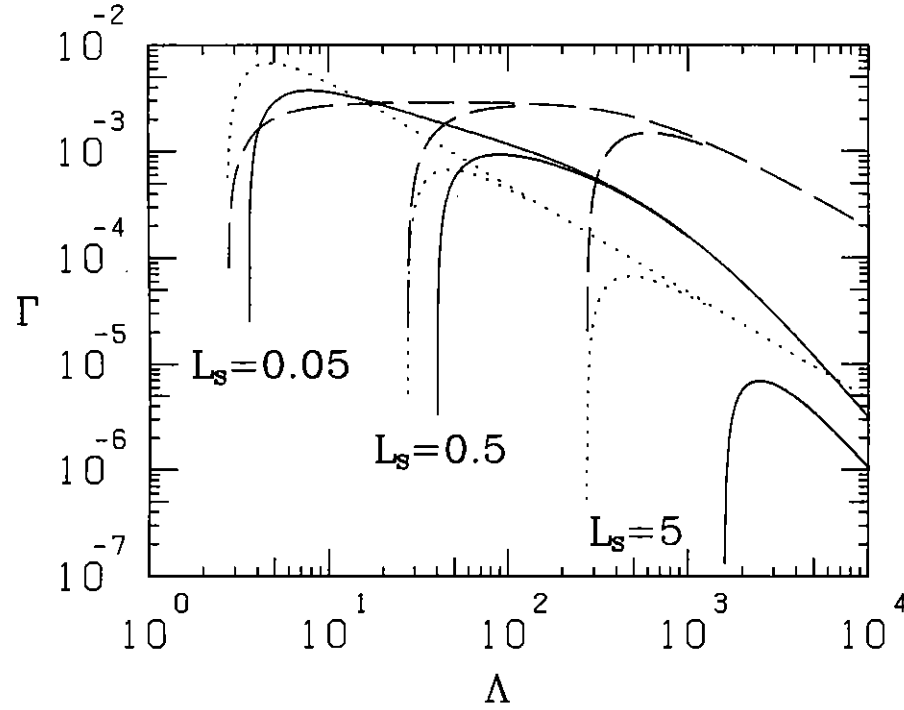


Fig. A5.6

Same as in Fig. 5.5 (with $\phi = 0$ and including f'_3 effects) but for $L_s = 0.05, 0.5$ and 5 . Solid lines are the results obtained from the present kinetic theory, dashed lines refer to the 2-fluid theory of Mochizuki et al (5.08) and dotted lines refer to the collisionless Weibel theory (5.16).

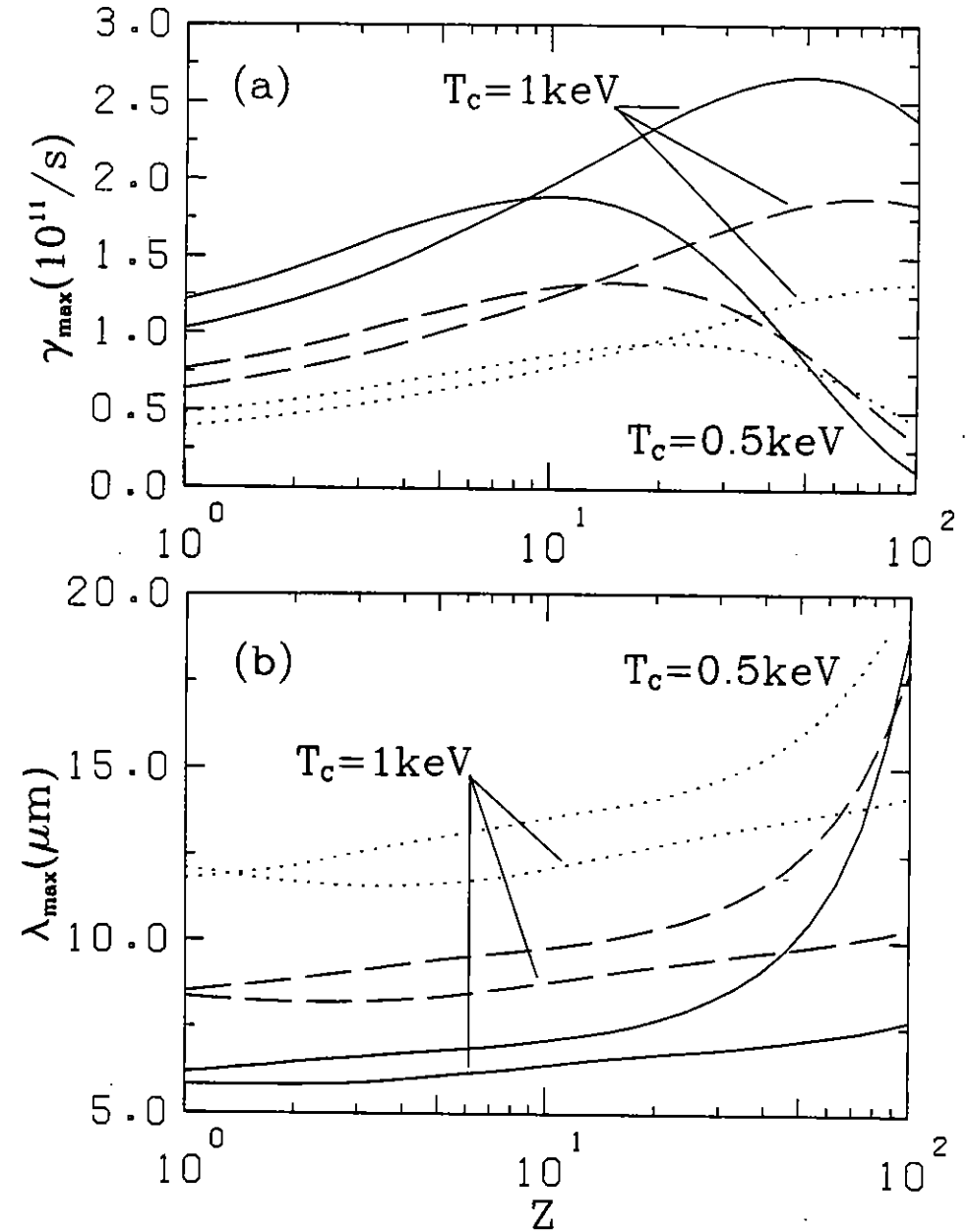


Fig. A5.7

Plot of: (a) maximum growth rate $\gamma_{\max}$ and (b) optimum wavelength $\lambda_{\max}$, as function of Z . Dotted lines are for $n_c = 5 \times 10^{21} \text{ cm}^{-3}$, dashed lines for $n_c = 10^{22} \text{ cm}^{-3}$ and solid lines for $n_c = 2 \times 10^{22} \text{ cm}^{-3}$. The equilibrium configuration is the same as in Fig. 5.5.

however, that these results do not allow for the n_e , T_e , Z dependence of the equilibrium. Since the instability depends on a high velocity moment of the distribution, namely the hot electron tail, one should strictly perturb an equilibrium derived from a non-linear Fokker-Planck simulation of the transport in a plasma (A5.14,A5.15).

In conclusion, it has been shown, by means of a novel kinetic model, that the collisional Weibel instability may explain the experimentally observed filamentary structures in laser-plasmas. This instability would appear to be most dangerous for high- Z targets.

I wish to thank the participants of the CECAM Workshop on 'Transport, Interaction and Instabilities in Laser-Fusion Plasmas', in October 1983, for many useful comments.

A5.2.4 Inverse Bremsstrahlung in the presence of Large Amplitude Langmuir Waves

R G Evans (RAL)

The absorption of an electromagnetic wave propagating in a fully ionised plasma is given by the imaginary part of the wave number as

$$k_i = \frac{1}{2} k_v \frac{\omega_p^2 \nu / \omega^3}{(1 - \omega_p^2 / \omega^2)^{3/2}}$$

where k_v is the vacuum wavenumber, ω and ω_p are the laser plasma frequencies and ν is the effective collision frequency. The equation is valid for $\nu \ll \omega$. For a thermal plasma with V_e defined by $mV_e^2 = kT_e$ the collision frequency is

$$\nu = \text{const.} \frac{Z^2 e^4 n_i}{m^2 V_e^3}$$

When the electromagnetic wave reaches a high enough intensity that the oscillating velocity $V_o = eE/m\omega > V_e$ the effective collision frequency is modified (A5.17,A5.18) by the distorted electron distribution function and if $V_o \gg V_e$ one has approximately

$$\nu = \text{const.} \frac{Z^2 e^4 n_i}{m^2 V_o^3}$$

Since the electromagnetic energy density $\epsilon \propto V_o^2$ one now has the seemingly paradoxical result that the absorbed energy decreases as the intensity increases.

A third situation arises if there is a driven large amplitude plasma wave whose oscillatory velocity V_{oe} satisfies $V_{oe} > V_o, V_e$. This situation arises for example in the "beat wave accelerator" (A5.19) where two electromagnetic pump waves drive a resonant plasma oscillation. At the saturated limit one has

$$c \approx V_{ol} \gg V_o \gg V_e$$

In this case the effective collision frequency

$$\nu = \text{const.} \frac{Z^2 e^4 n_i}{m^2 c^3}$$

and the absorption is much less than would be calculated from "normal" strong field theory. The Langmuir wave has a phase velocity close to c and itself loses energy by collisional damping.

As in the case of the reduction of strong field inverse bremsstrahlung with intensity, the energy density in the plasma wave exceeds that in the electromagnetic wave by $(V_{oe}/V_o)^2$ but the absorption coefficient is less by $(V_o/V_{oe})^3$. Overall less thermal energy is input to the plasma by collisional damping of the Langmuir wave than by collisional damping of the laser waves.

When numerical values are inserted for the beat wave example quoted in ref. A5.19 this decrease in absorption coefficient is vital in making the absorption length greater than the suggested plasma length.

The constraints in the above equations are not quite equal since the Coulomb logarithm changes as the electron distribution function is strongly distorted. The order of accuracy is

$$\frac{\ln V_o^2 / V_e^2}{\ln \Lambda} \text{ for the strong field calculations and } \frac{\ln V_{oe}^2 / V_e^2}{\ln \Lambda} \text{ for the strong}$$

Langmuir wave.

A5.2.5 Discrete Ordinates Calculations of Heat Flow in a Steadily Ablating Plasma

T H Kho (Imperial College)

An understanding of heat transport, by electrons, in the ablating region of laser-implosion targets, is important for the interpretation of experimental results and the design of targets for laser fusion. To gain insight into the effect of the ablation on the heat flow, a self-consistent kinetic calculation has been performed for the case of steady state ablation in planar geometry.

Considered is a fully ionised (Al^{13+}) plasma, steadily ablating from a cold to a hot boundary, in analogy to the ablation and critical surfaces respectively. The plasma carries no net current and the boundary conditions are that the inward-flowing part of the electron distribution functions at the hot and cold boundaries are half-Maxwellian with prescribed temperatures, T_h and T_c .

Electron transport is described by the one-dimensional transport equation

$$v \frac{\partial f}{\partial x} - eE \frac{\partial f}{\partial v} = \frac{\partial f}{\partial t} \quad \text{ie}$$

where f is the electron-distribution function in phase space, x the preferred direction, v the electron speed, E the electric field, and $(\partial f / \partial t)_{ie}$ the electron-ion scattering. This is solved in conjunction with the plasma equation of motion, $m_i n_i u (\partial u / \partial x) = - \partial P / \partial x$, where $P = m_e \int_{-\infty}^{\infty} v^2 f dv$. Here, m_e is the electronic mass. For simplicity, in the transport equation, electron self-collisions have been ignored. In treating the electron-ion scattering, the ions are assumed to be fixed. This is a reasonable approximation as the plasma flow velocity is of the order of the ion sound speed, $c_s(x)$, which is much smaller than the electron thermal speed $v_{te} (v_{te} / u \sim (m_i / m_e) < 1)$. Only subsonic flows are allowed; planar geometry cannot support steady supersonic flows. In the equation of motion, the ion

pressure is neglected because it is small compared with the electron pressure ($Z=13$, (ion temperature per electron temperature) < 1).

These equations are solved numerically using two nested iterations. An inner iteration solves the transport equation for an assumed electric field. Using the electron total energy $Q (Q = \frac{1}{2} m_e v^2 - e\phi(x))$, $\phi(x)$, is the electric potential, and its momentum perpendicular to the direction of the electric field $L (= m_e v(1-\mu^2)^{1/2}, \mu = v_x / v)$ as the independent variables, the transport equation can be expressed as

$$\frac{\partial f}{\partial x} = \frac{\partial f}{\partial t} \quad \text{Q.L. ie}$$

The equation above can be solved by integrating along trajectories of constant Q and L in phase space. This is done numerically by mapping these trajectories onto an angular mesh across the spatial domain.

The outer iteration searches for the self-consistent electric field to couple the electron and ion motion. The equation of motion which must be satisfied upon convergence is used to correct the electric field according to the formula

$$4\pi \frac{\partial}{\partial x} \int (Q + e\phi^*) \mu^2 f dv = - m_i n_i u \frac{\partial u}{\partial x}$$

Here $\phi^*(x)$ is the corrected potential to be used in the next iteration. All the other quantities are evaluated at the current iteration. Convergence is fastest with no ion flow (in which case the electron pressure is constant). The number of outer iterations needed for the simulations presented below ranged from 5 to 20.

Calculations were made with $T_h / T_c = 2$ for different mass fluxes and system thicknesses parameterized by $D / \bar{\lambda}$, where D is the spatial separation between the boundaries and $\bar{\lambda}$ is the mean free path of an electron at the hot boundary with kinetic energy $\frac{1}{2} k(T_h + T_c)$. Figure A5.8 illustrates a simulation with $D / \bar{\lambda} = 10$ in which the plasma flow attained a Mach number (flow speed per ion sound speed) of 0.74 at the hot boundary. The density, temperature (kinetic), and total electron pressure are shown.

Figure A5.9 shows the effect of increasing ablating flow on the electronic

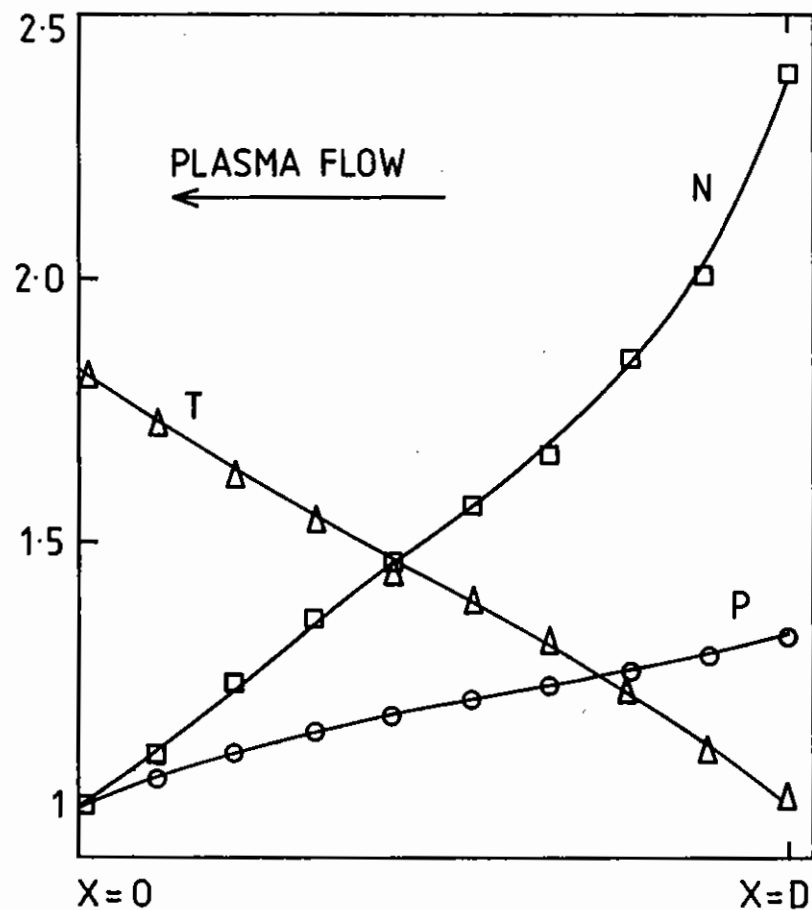


Fig. A5.8 The electron temperature (T), total electron pressure (P) and density (N) for a simulation with $T_h/T_c = 2$ and $D/\lambda = 10$. The hot boundary is at $X = 0$ and the cold boundary, $X = D$. The units are arbitrary.

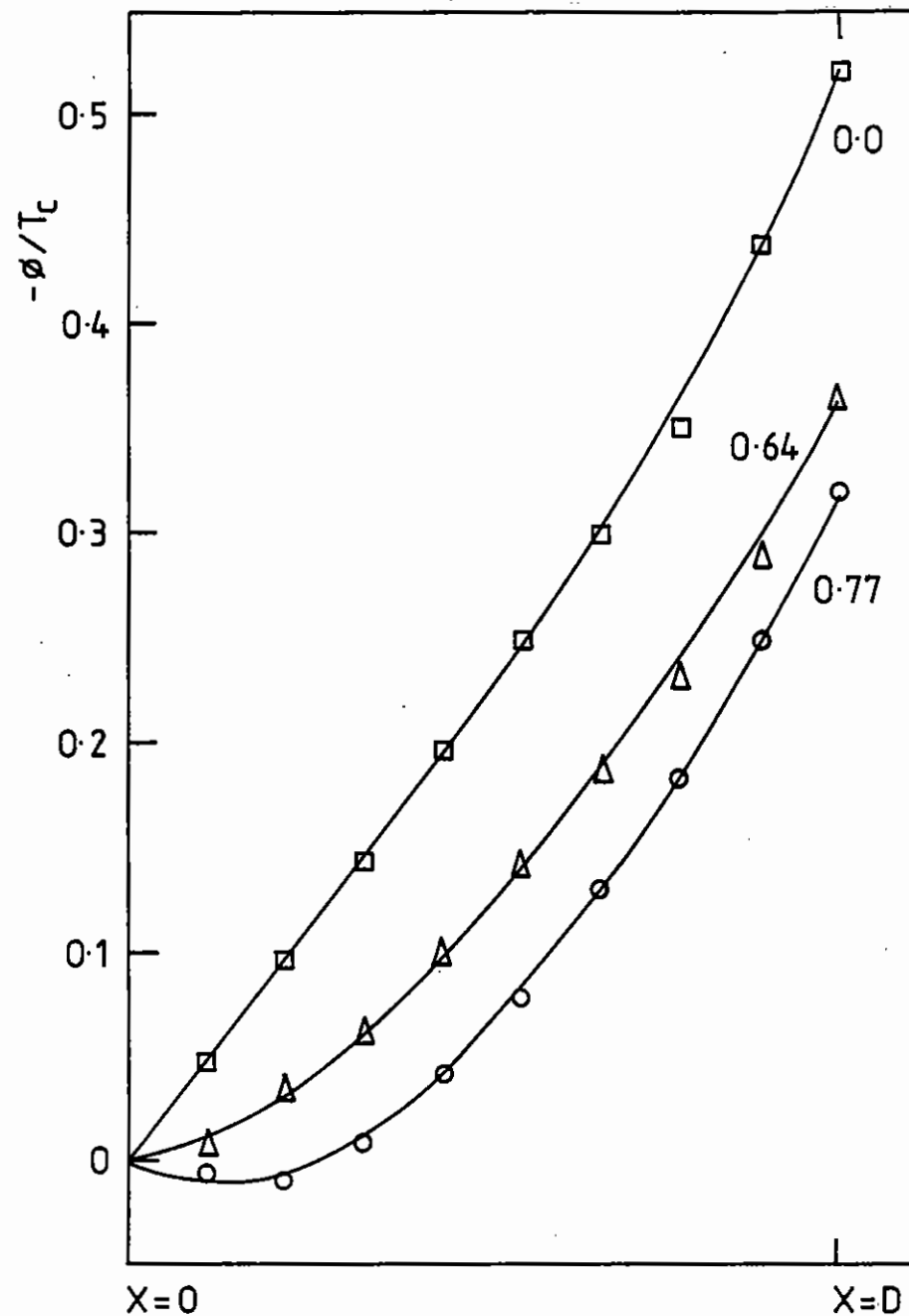


Fig. A5.9 Effects of increasing steady state ablation on the electron potential ($-\phi(x)$) along a system with $D/\lambda = 10$. Each curve represents a simulation with a given (constant) mass flux (m, n, u) and is labelled by the Mach number the plasma flow attained at the hot boundary ($X = 0$). The top curve represents a simulation with no ion motion ($\partial P / \partial x = 0$).

potential $-\epsilon\phi(x)$. For sufficiently high mass flow the potential becomes non-monotonic near the hot boundary.

In these simulations we have ignored the existence of electrons with total energies less than the escape energy of the potential trough and which could, hence, be trapped in it (the total energy is a constant of motion here because of the Lorentz plasma approximation). It is possible that in experiments with high-Z targets some electrons are trapped just behind the critical surface in this way.

Figure A5.10 shows the normalized heat flow Q/Q_{SP} (calculated center-of-mass heat flow per Spitzer-Harm thermal conduction evaluated at the local values of (λ/L_T) : λ is the mean free path of an electron with energy $kT(x)$ and L_T is the temperature scale length). Each bold curve represents Q/Q_{SP} measured along the temperature front. The dotted curves represent simulations with ion motion and are labelled by the Mach number the plasma flow attained at the hot boundary. From left to right, each pair of bold and dotted curves corresponds to systems with thicknesses $D/\bar{\lambda}=20, 10, 5$ and 0.5 . For all cases, the lowest point of each curve corresponds to that spatial point nearest the hot boundary (critical surface), i.e., at the top of the temperature front.

It can be seen that inclusion of steady-state ablation tends to increase Q/Q_{SP} . As the system gets less collisional the difference between the results with and without ion motion decreases. For both cases, in the collisionless limit all macroscopic quantities (including the potential) are constant. The difference between the two cases arises then from the fact that in the fluid frame the flow of the electrons from the hot boundary is shifted inwards with ion motion. For the subsonic flows here, the difference is quite small (Fig.A5.10, $L_T/\lambda \rightarrow 0$). If there is no ion motion then the electric field in the direction of the heat flow increases as the system becomes more collisional. The inclusion of ion motion does two things, first it reduces this electric field as described above and second it raises the electron pressure gradient to drive it.

The increased pressure gradient is effected by a weak smoothing of the

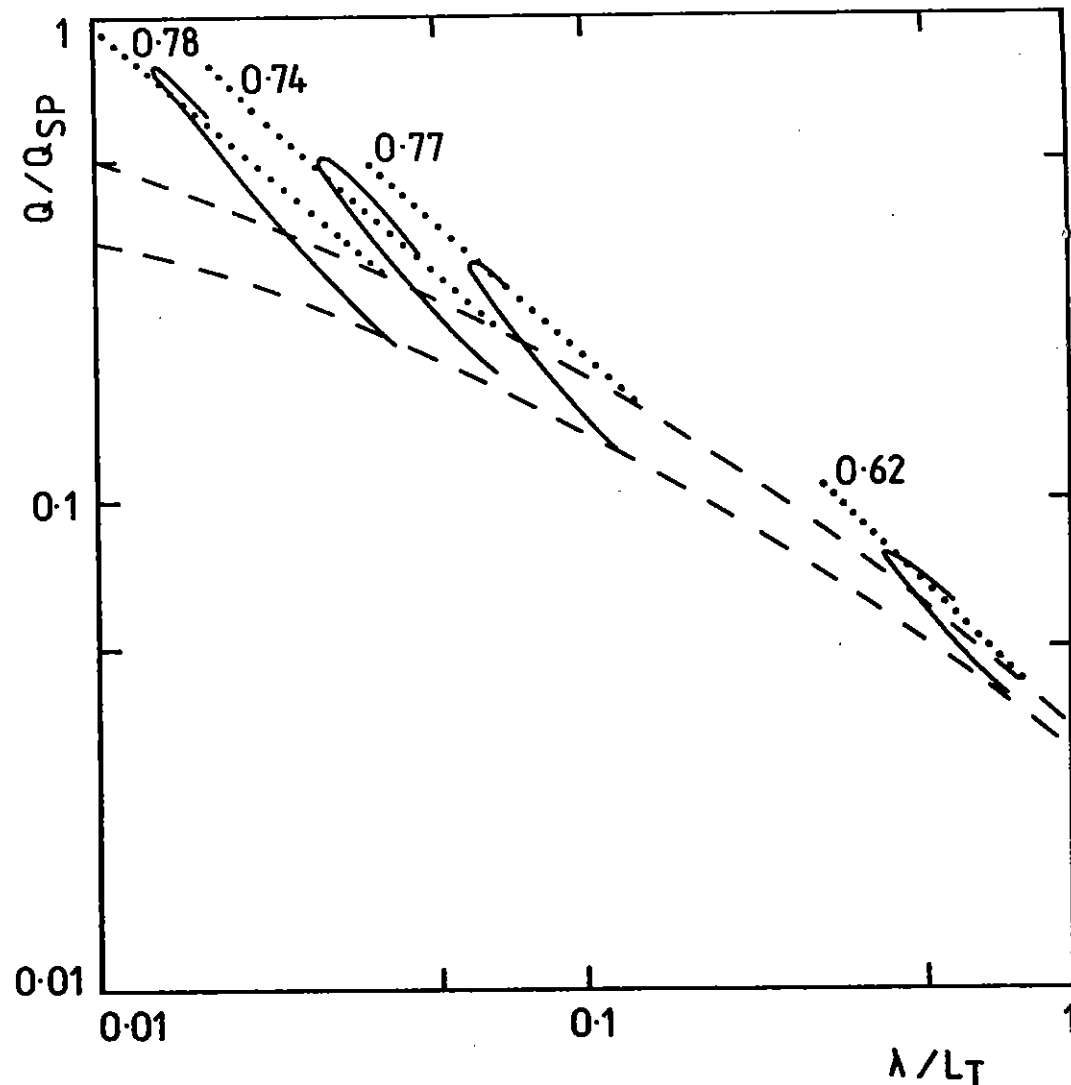


Fig. A5.10

Comparison of Q/Q_{SP} (calculated heat flow/Spitzer-Harm thermal conduction for a Lorentz plasma evaluated at the local values of λ/L_T). Each solid curve represents Q/Q_{SP} along a system of given D/λ with no ion motion. The effect of including a steady state ablation in that system is represented by the dotted curve immediately above the bold curve. The dotted curves are labelled by the Mach number the plasma flow attained at the hot boundary ($X = 0$). The dashed curves approximate the heat flows at the top of the temperature fronts.

temperature profile and a more marked increase in density near the cold boundary. The increased density makes the system more collisional and would tend to lower the heat flow but this effect is not enough to compensate for the lower electric field.

As the system gets more collisional, the difference in the electric field with and without steady-state ablation becomes more important. This is reflected in the difference in the heat flows between the two cases with increasing collisionality.

However, this difference is still small. The ion motion, for the case studied here, is not found to have any major influence on the heat flow.

A5.2.6 Non-Spitzer Heat Flow in a Laser-Produced Spherically-Symmetric Steadily-Ablating Plasma

A R Bell (RAL)

It is well known that thermal transport in laser-produced plasmas is not consistent with the linear theory of conductivity developed by Spitzer and Harm. The large heat flow requires non-linear solution of the Vlasov-Fokker-Planck (VFP) equation for electrons. The details of heat flow in any situation are sensitive to the environment in which it occurs and it is important to include the hydrodynamics of ablation when considering heat flow in laser-produced plasmas. In last year's annual report (A5.20), calculations of heat flow in a steadily ablating plasma in planar geometry were presented. In that case, it was found that, although the conductivity differed from the Spitzer conductivity where the heat was input, the difference was not sufficient to make a measurable difference to the temperature profiles. The conclusion of that work was that in a steadily ablating plasma, the Spitzer conductivity gives a good description of heat flow. However, the special case of planar geometry is applicable only to experiments at low laser irradiance in which the diameter of the laser focal spot, or the target diameter, is much larger than the distance between the solid target and the absorption surface. Deviation from the Spitzer conductivity has been experimentally measured predominantly at high laser

intensities (A5.23) for which the planar approximation is not valid. Under such circumstances, the plasma flow diverges as it expands away from the solid. This leads to cooling in the underdense plasma. Consequently, the temperature is at a maximum at the absorption surface and heat is conducted away in both directions, towards the solid target and also into the underdense plasma. A further difference from the planar case is that the plasma flows through a sonic surface (Mach number equal to one) in the region between the solid and absorption surfaces, thus allowing supersonic flow near the absorption surface. This results in a larger heat flow relative to the free-streaming limit,

$Q_f = n_e kT (kT/m_e)^{1/2}$, and steeper temperature gradients relative to the electron mean free path than is possible in the planar case. Because of these important differences between planar and divergent flows, and the tendency for non-Spitzer heat flow to be measured experimentally when the heat flow is strongly divergent, it is important to include divergent flow in numerical modelling. Here we extend our previous calculations into spherical geometry. Apart from the changes to the spatial terms needed by spherical geometry, the only major difference in the equations concerns the collision term. Because the conditions are more extreme in spherical geometry, the electron distribution deviates more strongly from Maxwellian, and we remove some of approximations made to the collision term in the planar calculation. The collision coefficients used in the spherical calculation are

$$C_{2+1} = \frac{18}{Z+1} \frac{1}{\lambda} \frac{1}{v^2} \frac{n}{v_0^2}$$

$$D_1 = \frac{18}{Z+1} \frac{1}{\lambda_0} \frac{1}{v^3} \left(\frac{n k T}{m_e} - \int \frac{4\pi}{3} (w^4 - v^3 w) f_0(w) dw \right) \left[1 - \int \frac{4\pi w^2}{v} f_0(w) dw / n \right]^{-1}$$

$$D_2 = \frac{9n}{\lambda_0 v}$$

where the symbols have the same definition as in last year's annual report. The coefficients C and D_2 are the same as before, but D_1 is weighted such that the ratio D_1/C is correct even though the magnitude of C and D_1 is not correct at low electron velocities. Because the collision term dominates

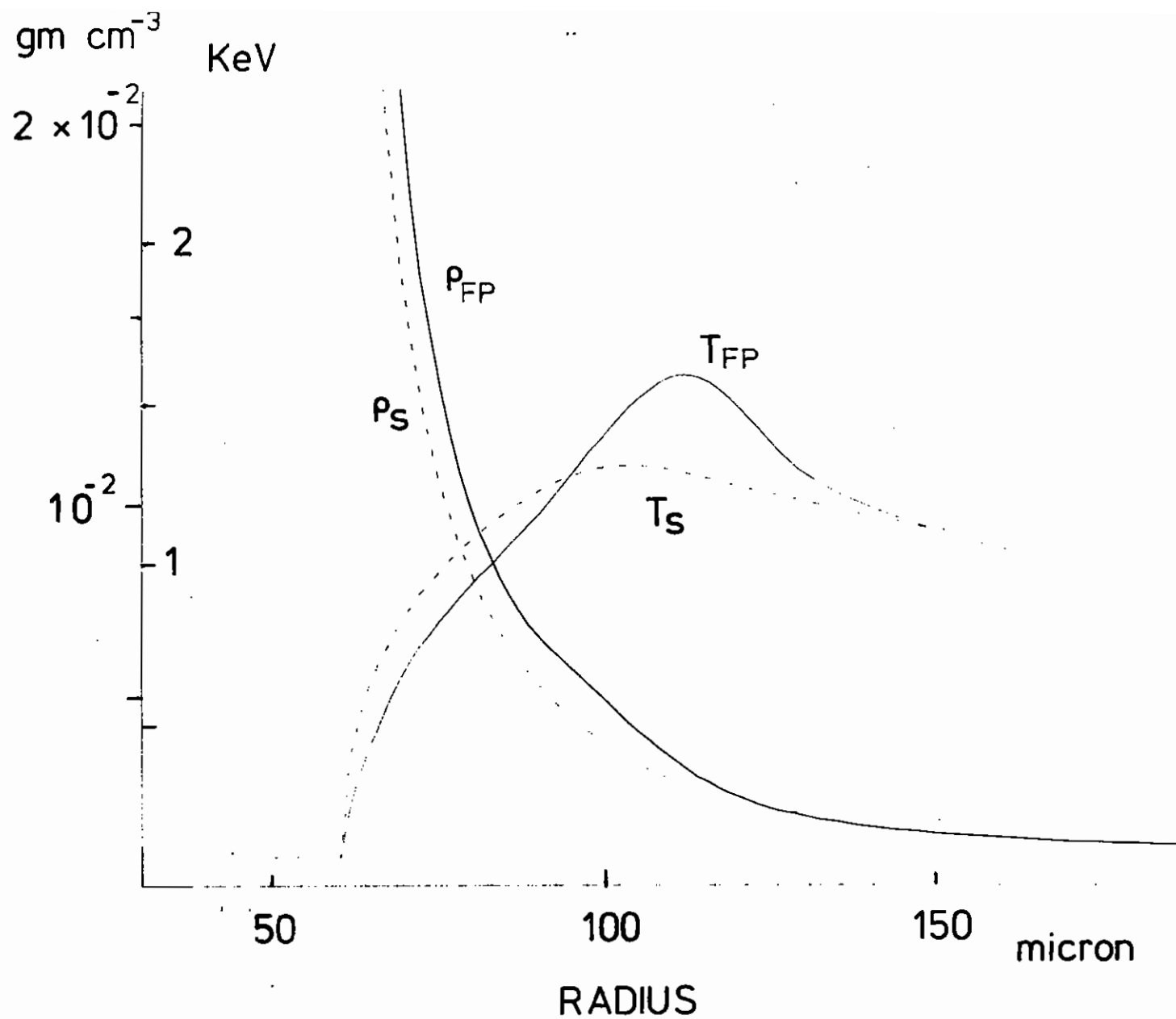


Fig. A5.11 Temperature (T) and density (ρ) as functions of position. The absorbed irradiance is $1.7 \times 10^{14} \text{ W cm}^{-2}$. The full lines are the temperature and density profiles found by self-consistent solution of the hydrodynamic equations with the VFP equation. The dashed lines are the profiles found by solution of the hydrodynamic equations with the Spitzer conductivity instead of the VFP equation. The electron-ion mass ratio is $m_i/2m_e = 3672$, and $Z = 15$.

the spatial term at low velocity it is the relative magnitude of these coefficients, as opposed to their absolute magnitude, which determines the electron distribution. The hydrodynamics of the bulk plasma flow are defined by conservation of mass, momentum and energy as described by equations 7 to 9 in reference A5.21, except that the laser energy is deposited over a volume of plasma instead of at an infinitesimally thin surface. The spatial distribution of energy absorption is arbitrarily imposed as an input function to the calculation. The calculation is made steady-state by including a mass source at zero radius. As before, the ions are assumed to be cold. The major additional difficulty introduced by spherical geometry is the integration of the hydrodynamic equations through the sonic surface. This has been discussed by Gitomer, Morse and Newberger (A5.22), and their approach of integrating away from the sonic surface in both directions has been adopted here. Although the equations can be placed in a dimensionless form as in the planar case, the solutions in spherical geometry consist of a family of solutions instead of a single solution. This is caused by the introduction of an extra variable, the radius of the target. Each member of the family of solutions can be characterised by the ratio of the stand-off distance of the absorption surface from the solid target to the target radius. This means that, for a given target radius, a different solution is required for each value of the laser irradiance since the stand-off distance increases with irradiance.

We give results for an absorbed irradiance of $1.7 \times 10^{14} \text{ W cm}^{-2}$ at a laser wavelength of $1 \mu\text{m}$ and a target diameter of $120 \mu\text{m}$. The temperature and density profiles are shown in Fig. A5.11. The profiles outside the solid surface are defined by self-consistent solution of the VFP electron equation and the fluid equations for conservation of mass, momentum and energy. The steady state approximation leads to non-uniform density and temperature within the solid and we replace the mathematically correct steady state solution with a more realistic uniform density and temperature. As a consequence, this calculation cannot give information about low level preheat at temperatures less than that imposed in the solid, which is 0.06 times the maximum temperature in the system. The temperature and density profiles given by the dashed lines in Fig. A5.11 are calculated by solving the same set of equations except that the heat flow is calculated from the

Spitzer conductivity instead of the VFP equation. In contrast to the planar results, there is now a significant difference between the Spitzer and VFP temperature profiles. The more notable differences are that the temperature maximum is higher, the shape of the temperature front is less convex and the density gradient is larger in the region where energy is absorbed. A surprising result is that the stand-off distance of the absorption surface is larger in the VFP calculation. This contrasts with the result given by 'flux-limited' theories of heat flow and is caused by an enhancement in the conductivity at the base of the heat front resulting from relatively high energy electrons with long mean free paths, originating higher up the temperature front, flowing freely into the cold plasma. This could be regarded as a preheat effect, although its effect is only to reduce the convexity of the heat front instead of making it actually concave. The difference between the planar and spherical calculations (i.e. low and high irradiance) implies that the phenomenon of 'flux inhibition' should only be apparent at high irradiance unless there are other factors reducing the heat flow. Calculations at other irradiances would be needed to calculate the threshold for inhibition. At present it is only possible to state that it occurs at an irradiance below about $10^{14} \text{ W cm}^{-2}$. It is interesting to note that, in Rosen's recent review (A5.23) of experimental results, he comes to the conclusion that there is a threshold for non-Spitzer heat flow at about $3 \times 10^{13} \text{ W cm}^{-2}$. Heat flow is frequently investigated experimentally by measuring the mass ablation rate, and comparing this with numerical hydrodynamic simulations which model heat flow by imposing an upper limit of FQ_f on the heat flow carried by a plasma. The value of the 'flux limit' F is varied until a good fit with experiment is obtained. In this approximate flux-limited treatment of conductivity the mass ablation rate decreases as F is reduced i.e. as the reduction from the Spitzer conductivity is made stronger. At the irradiance of $1.7 \times 10^{14} \text{ W cm}^{-2}$, we calculate that the mass ablation rate derived from the VFP equation is very close to the value given by the Spitzer calculation, and is indeed slightly higher, although only by 3 percent. However, it must be remembered that the VFP calculation is in steady state. Because the conductivity is in fact reduced below the Spitzer value over the greater part of the temperature front, a time-dependent VFP calculation of the initial setting up of the temperature front might be expected to show the front still penetrating into the solid more slowly than

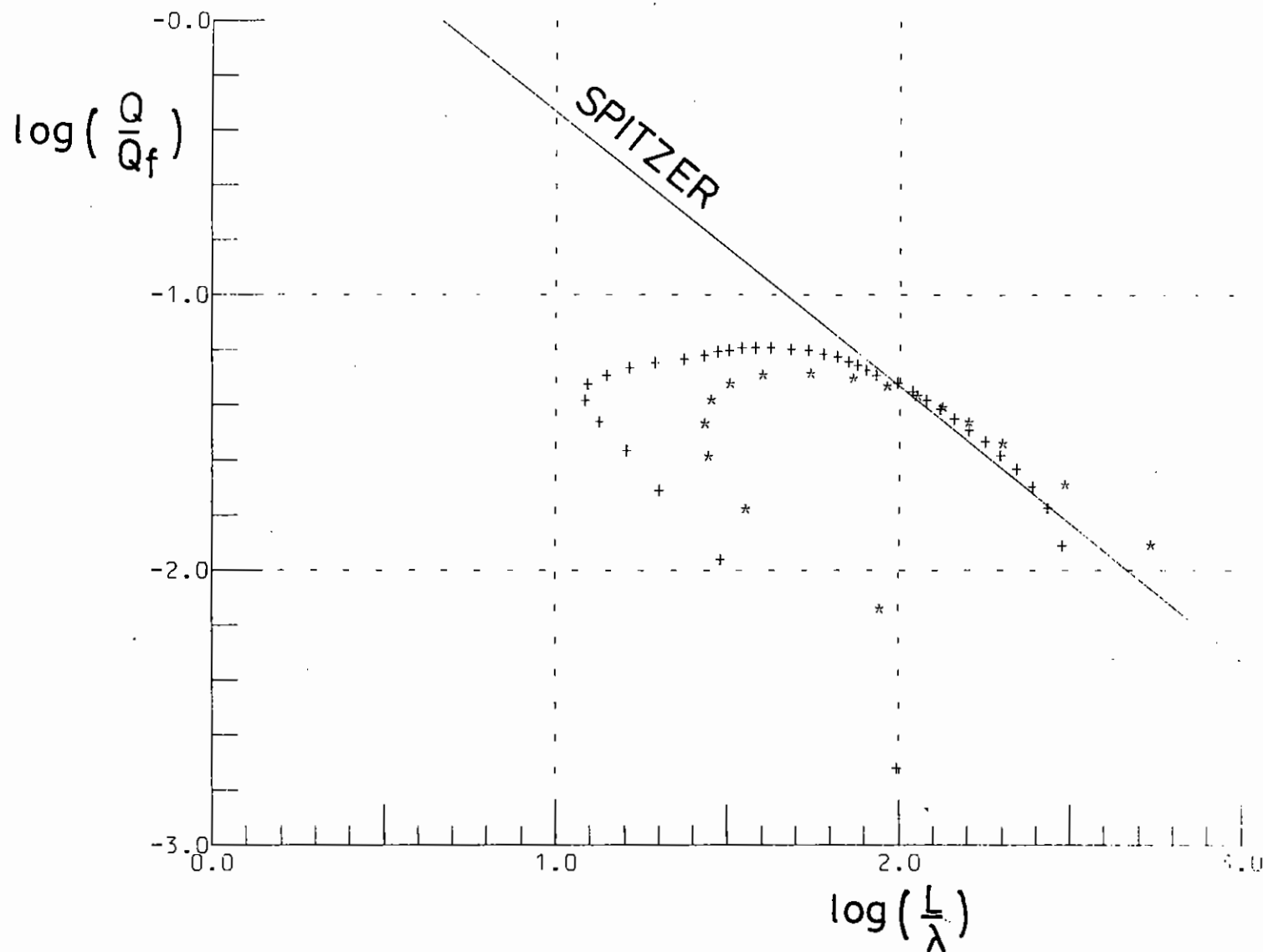


Fig. A5.12 $\log(Q/Q_f)$ plotted against $\log(L/\lambda)$ where $Q_f = n_e kT(kT/M)^{1/2}$, $L = T/\nabla T$ and λ is the mean free path of an electron with energy $3kT/2$. The irradiance is $1.7 \times 10^{14} \text{ W cm}^{-2}$. All quantities are calculated locally. Asterisks represent inward heat flow, crosses outward heat flow. The solid line is the Spitzer conductivity. The electron-ion mass ratio is $m_i/Zm_e = 3672$, and $Z = 15$.

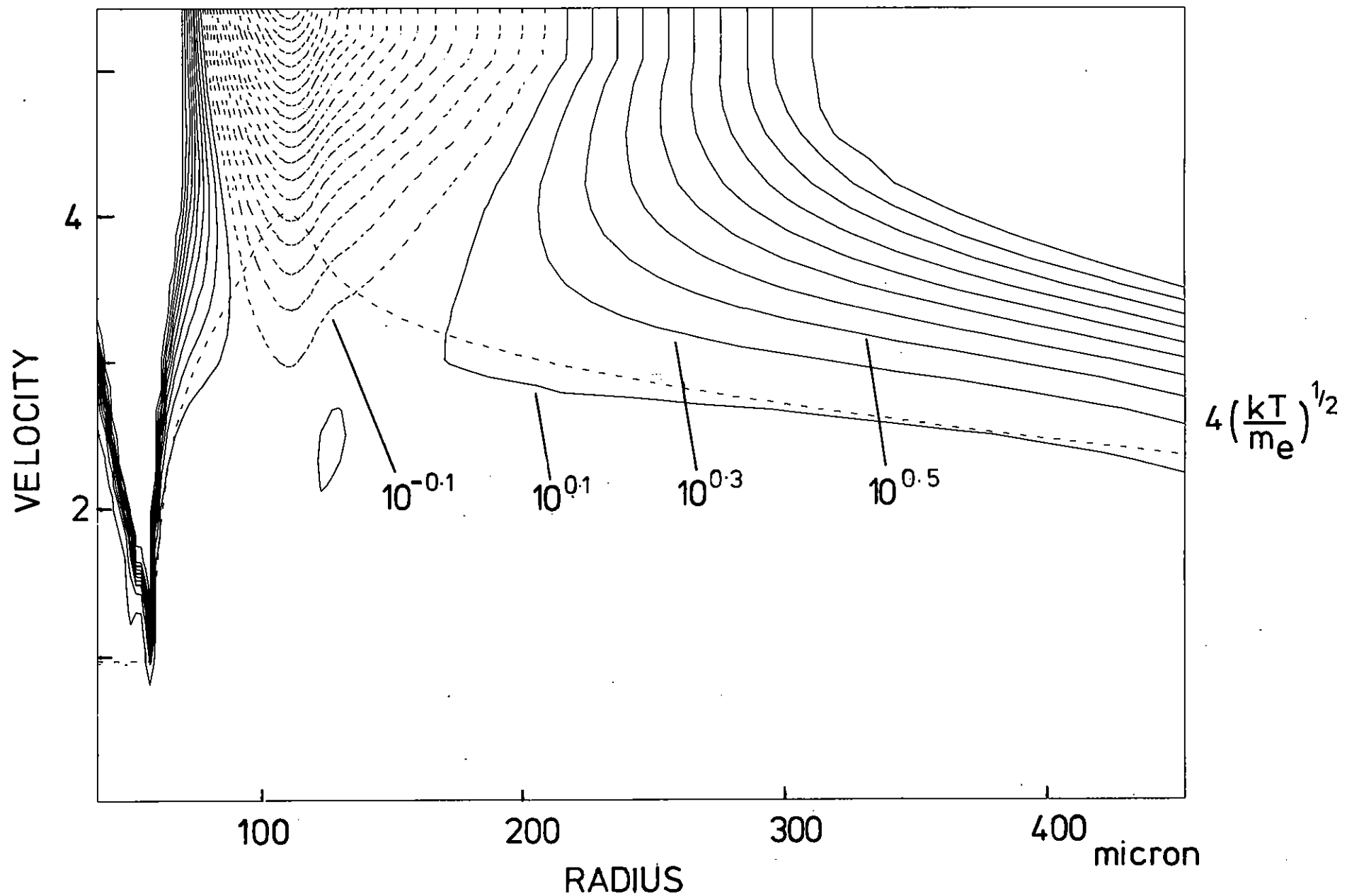


Fig. A5.13 Contour plot in (x, v) of the electron density in phase space relative to the Maxwellian density $f_0(v) \exp(mv^2/2kT)/f_0(0)$. The absorbed irradiance is $1.7 \times 10^{14} \text{ W cm}^{-2}$. The contours are spaced logarithmically at $10^{-0.1}$, $10^{0.1}$, $10^{0.3}$, $10^{0.5}$... $10^{1.9}$. Negative contours are dotted. The velocity at which electrons have four times their thermal velocity is given by the dashed line which is defined by $v = 4(kT/m_e)^{1/2}$. $m_i/Zm_e = 3672$ and $Z = 15$.

given by the Spitzer calculation and the mass ablation rate correspondingly reduced below the Spitzer rate. Our calculations predict that as laser pulse lengths are increased and experiments progress further into the steady-state regime, the mass ablation rate will agree more closely with the Spitzer value and will cease to be a good indicator of the level of heat flow.

We now consider the details of the heat flow calculated from the VFP equation and the explanation of its large departure from the Spitzer value. Figure A5.12 is a plot of the heat flow at each point on the computational grid as a fraction of the local free-streaming limit Q_f versus the local scale length of the temperature gradient $T/|\nabla T|$ in units of the local electron mean free path. The effective conductivity deviates significantly from the Spitzer conductivity which is given by the full line. The heat flow outside the absorption surface towards large radii (marked by crosses) reaches a maximum of around 6 percent of the free-streaming limit. The heat flow inwards towards the solid, marked by asterisks, reaches a maximum of 5 percent of the free-streaming limit while still deviating significantly from the Spitzer value. The points which lie above the Spitzer line represent heat flow at the base of the temperature gradients and constitute an enhanced conductivity. VFP calculations have not previously shown non-Spitzer heat flow at the low level of 5 percent of the free-streaming limit and we try to identify the mechanism which gives this result. Figure A5.13 is a contour plot of the phase space density of the isotropic part f_0 of the electron distribution function as a fraction of what it would be if the distribution were Maxwellian everywhere. At low electron velocities the distribution is close to Maxwellian, and deviates from Maxwellian predominantly at velocities greater than three times the local thermal velocity $(kT/m_e)^{1/2}$. Because the high-velocity electrons have long mean free paths and high mobilities the tail of the distribution is depleted where the temperature is high and enhanced where the temperature is low as high velocity electrons flow freely down the temperature gradient. This reduces the gradient in the electron phase space density below its value for Maxwellian distributions. This in turn reduces the diffusive electron flow and the heat flow is correspondingly reduced. This effect was present

in the planar calculation but is much larger in the spherical calculation for the following reasons: 1) In planar geometry the stand-off distance of the absorption surface from the solid is about 60 times the mean free path of an electron with energy $3kT_h/2$ where T_h is temperature at the absorption surface. In spherical geometry the existence of an extra characteristic length, i.e. the target radius, which is relatively short forces the stand-off distance to a smaller value. Thus, the gradients are shorter relative to the electron mean free path allowing non-local and non-linear effects to become more important. 2) Because the temperature falls away in both directions from the absorption surface, there are two directions in which electrons on the tail of the distribution at the absorption surface can escape compared with the one direction, towards the solid, in the planar calculation. Furthermore, the density decreases into the underdense plasma making the mean free path relatively long and allowing electrons to escape easily from the absorption surface to large radii. This increases the depletion of the high-velocity tail at the absorption surface. Consequently, the gradient in electron phase space density in the region between the absorption surface and the solid is reduced, reducing the heat flow to a greater degree in spherical geometry than in planar geometry.

A5.3 Hydrodynamics and Compression

A5.3.1 Rayleigh-Taylor and Kelvin-Helmholtz instabilities by vortex sheet method

S Jorna (St. Andrews)

The nonlinear development of the Rayleigh-Taylor and Kelvin-Helmholtz instabilities of the interface between two immiscible, incompressible, and irrotational fluids can be studied in two dimensions by representing the interface as a vortex sheet. The expression for the complex potential of a periodic distribution of discrete vortices dates back to Kelvin. More than fifty years ago, Rosenhead adopted this expression in a numerical study of the shape of a breaking wave. Only recently has the method yielded spike extensions in the Rayleigh-Taylor instability as large as twice the initial amplitude. Pullin's (A5.24) is arguably the clearest exposition, and in

what follows we adopt his notation.

For a sharp, periodic interface the formulation is exact for irrotational and incompressible fluids, and the numerical results might usefully be compared to those obtained from more conventional fluid codes.

Subject to the assumptions already mentioned, the complex potential at a point $z(x+iy)$ due to a row of vortices of strength κ spaced λ apart is

$$W(Z) = \frac{\kappa}{2\pi i} \ln \left(\sin \frac{\pi}{\lambda} (z - z_0) \right), \quad (1)$$

for a vortex at z_0 .

If the vorticity is distributed continuously, the potential for a vortex sheet is (we assume the constant, uniform convection velocity of the fluids is zero)

$$W(Z, t) = \phi(x, y) + i\psi(x, y) \\ = \frac{1}{2\pi i} \int_0^{\lambda} \ln \left(\sin (Z - z(a, t)) \right) \frac{\partial \Gamma(a', t)}{\partial a'} da', \quad (2)$$

where a is the Lagrangian position and where Γ is the total circulation defined as

$$\Gamma = \int \bar{v} \cdot d\bar{x} \quad (3)$$

so that, if we define $\bar{v} = \nabla \phi$, Γ is the change in ϕ around the circuit.

The position of the interface is identified by marker particles whose velocity has the general form (cf. Pullin (A5.24))

$$\dot{z} = \left(\frac{dW}{dz} \right)_p^* + \frac{1}{2} f(a) \frac{\partial \Gamma / \partial a}{(\partial z / \partial a)^*}. \quad (4)$$

$f(a)$ is an arbitrary function of the position a , and $(dW/dz)_p$ indicates that the principal value is to be taken.

The time-dependence of Γ follows from the dynamic condition that the pressure be continuous at the interface (ignoring the surface tension here does not significantly affect the results). It follows then from

Bernoulli's equation that (we set $g=1$)

$$\frac{\partial \Gamma}{\partial t} + \frac{1}{2} \left(\left| \frac{dW}{dz} \right|_1^2 - \left| \frac{dW}{dz} \right|_2^2 \right) \\ = -A \left(\frac{\partial}{\partial t} (\phi_1 + \phi_2) + \frac{1}{2} \left(\left| \frac{dW}{dz} \right|_1^2 + \left| \frac{dW}{dz} \right|_2^2 + 2y \right) \right), \quad (5)$$

where the subscripts 1 and 2 refer to the fluids above and below the interface, respectively, and

$$A = \frac{\rho_1 - \rho_2}{\rho_1 + \rho_2}. \quad (6)$$

The velocities of the marker particles can be calculated from definition (Eq.4) and the z -derivative of Eq.2 for an initial distribution. The arbitrary function $f(a)$ has been set equal to a constant, and to retain accuracy at large deformations marker particles are added - while conserving total vorticity - in regions where the radius of curvature is small.

The evolution of the spike and bubble for the pure Rayleigh-Taylor instability ($A=1$) is shown in Figure A5.14. The time is scaled by g . Departure from the linear growth ($\propto \cosh(2\pi)^{1/2}t$) is already significant at time 0.425. After time $t=1$ the tip of the spike is in free fall. Our calculations cease to give meaningful results after the amplitude has grown to roughly twice the initial amplitude; apparently a common fate.

For $A \neq 1$, the Kelvin-Helmholtz instability is also excited as is shown in Figure A5.15 for $A = 1/4$. It will be noted that the Kelvin-Helmholtz instability reduces the growth of the Rayleigh-Taylor spike.

5.3.2 Radiation Loss Driven Instabilities in Laser Heated Plasmas

R G Evans (RAL)

Over the last few years a number of observations have been reported (A5.7, A5.8, A5.25) of significant non-uniformities in the low density plasma

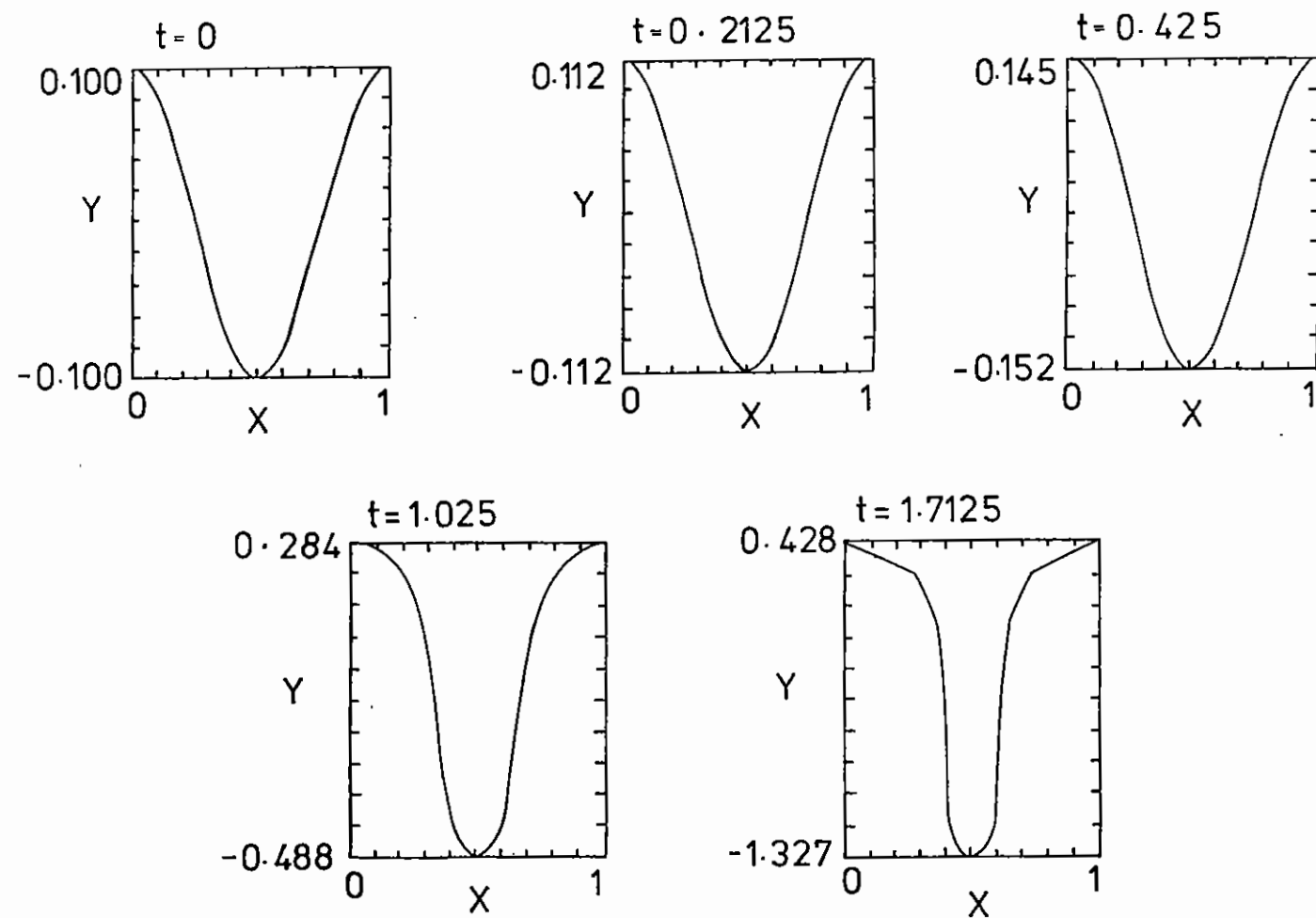


Fig. A5.14 Development of pure Rayleigh-Taylor instability for scaled times 0, 0.2125, 0.425, 1.025 and 1.7125.

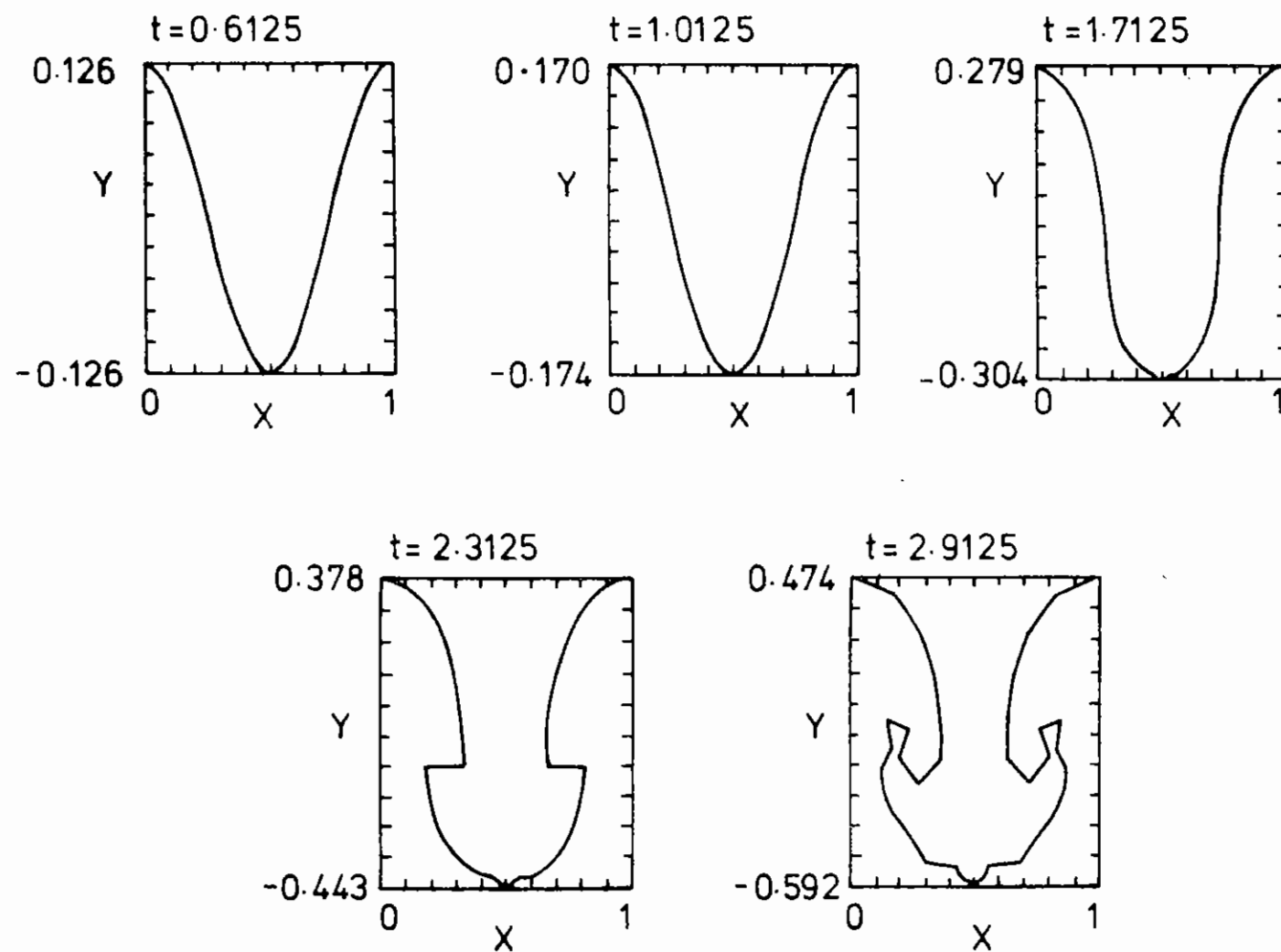


Fig. A5.15 Rayleigh-Taylor and Kelvin-Helmholtz instabilities ($A = 1/4$) for scaled times 0.6125, 1.0125, 1.7125, 2.3125 and 2.9125.

ablated from a solid surface by a high intensity laser beam. In these experiments the target is not generally irradiated uniformly but it would be expected that non-uniformity of ablation near the solid material would be symmetrised as the material tends towards a spherical expansion into the vacuum. The observation that the non-uniformities persist far out into the ablation flow has caused people to suggest that plasma instabilities may be present.

Instabilities driven by various velocity anisotropies have been suggested including a magneto-thermal heat flow instability(A5.10) and a Weibel-like instability (A5.26). Self-focussing of the laser light and filamentation are probably excluded since the plasma "jets" are not parallel to the incident laser light and persist for some time after the laser pulse(A5.7). The observed non-uniformities are markedly Z dependent, being virtually non-existent on plastic (CH) targets, observable on Al or SiO₂ targets and quite dramatic on targets of Au or W. This Z dependence is consistent with the magneto-thermal and Weibel instabilities through the Z dependence of the thermal conductivity but has also produced the suggestion that the observations could be explained by optically thin radiation losses (A5.9). The radiative loss driven instability has been analysed for a rather artificial plasma equilibrium and shown to produce qualitatively the correct scale lengths and growth rates. Recently an attempt to incorporate radiative loss terms in an analytic steady ablation model (A5.27) has suggested that this instability is stabilised by convection.

The simple equilibrium analysed in ref.A5.9 is not really applicable to a laser irradiated plasma where the convective flow of energy and thermal conduction play a major part in defining the quasi-steady state. The most straightforward, but admittedly not the most elegant solution is to use a two dimensional hydrodynamic computer model which solves the full time dependent problem. The code used here is named POLLUX (A5.29) and solves the fluid equations on a sliding Eulerian mesh with (r,z) cylindrical symmetry. The code includes heat flow, laser energy deposition, an optically thin radiative loss function and separate electron and ion radiative loss function used is essentially that calculated by Summers and McWhirter (A5.30) for coronal equilibrium but is multiplied arbitrarily by

the exponential factor $\exp(-T_e/200\text{eV})$ to reduce the radiative losses in the optically thick part of the spectrum below $h\nu=600\text{eV}$. The radiative loss is included in the electron energy balance equation and no attempt is made to calculate the radiation transport.

The starting conditions for a typical computer run are shown in Fig.A5.16. The computational mesh has 60 x 30 mesh points covering 100 μm x 50 μm . The laser beam is incident from the right and has a wavelength of 0.53 μm , and a mean intensity of 10^{14}Wcm^{-2} . In order to provide an initial seed to any instability the laser beam is sinusoidally modulated with an intensity ratio 2:1, with one complete cycle in the r-direction; this is illustrated schematically on the right hand side of Fig.A5.16. The target was usually a thin foil of material 5 μm thick, purely for the reason that these simulations use less computing time than for a solid target. One run was repeated for a solid target to show that target acceleration does not affect the conclusions.

Simulation results for the iso-density contours on a plastic(CH) target are shown in Fig.A5.17 both with and without the radiative loss term in the energy equations. The contour plot at 300ps shows that the foil has moved farthest where the laser intensity was greatest but at 340ps the curvature of the one half critical density, $n_c/2$, surface has reversed due to the increased expansion of the underdense plasma in the hotter region. When the radiative loss terms are included the result at 300ps shows a noticeable perturbation of the $n_c/2$ surface but this does not persist through to 340ps. Experimentally no "jets" are observed on plastic targets(A5.7).

With a gold target the density contours in the absence of radiative losses show the same qualitative behaviour as for plastic, Fig.A5.18. When the radiative loss terms are included in the simulation the results are dramatically different as shown in the lower part of Fig.A5.18. High spatial frequency ripples are observed at 300ps, leading at 340ps to a pronounced axial jet about 20 μm long and 10 μm radius. The jet is confined to plasma densities below n_c in contrast to the Rayleigh-Taylor accelerational instability which is dominant at densities close to the solid density.

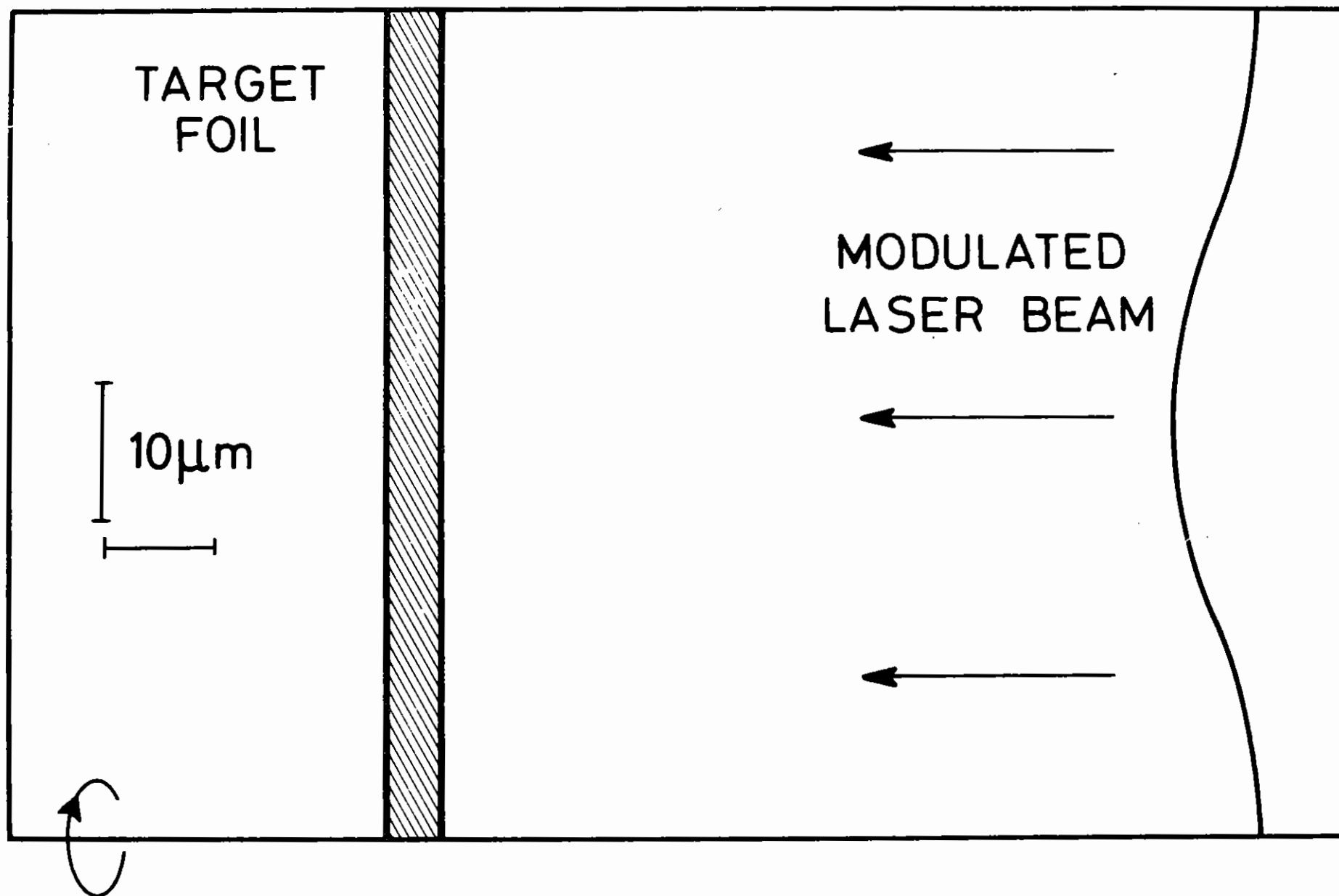


Fig. A5.16 Initial conditions for a typical computer simulation.

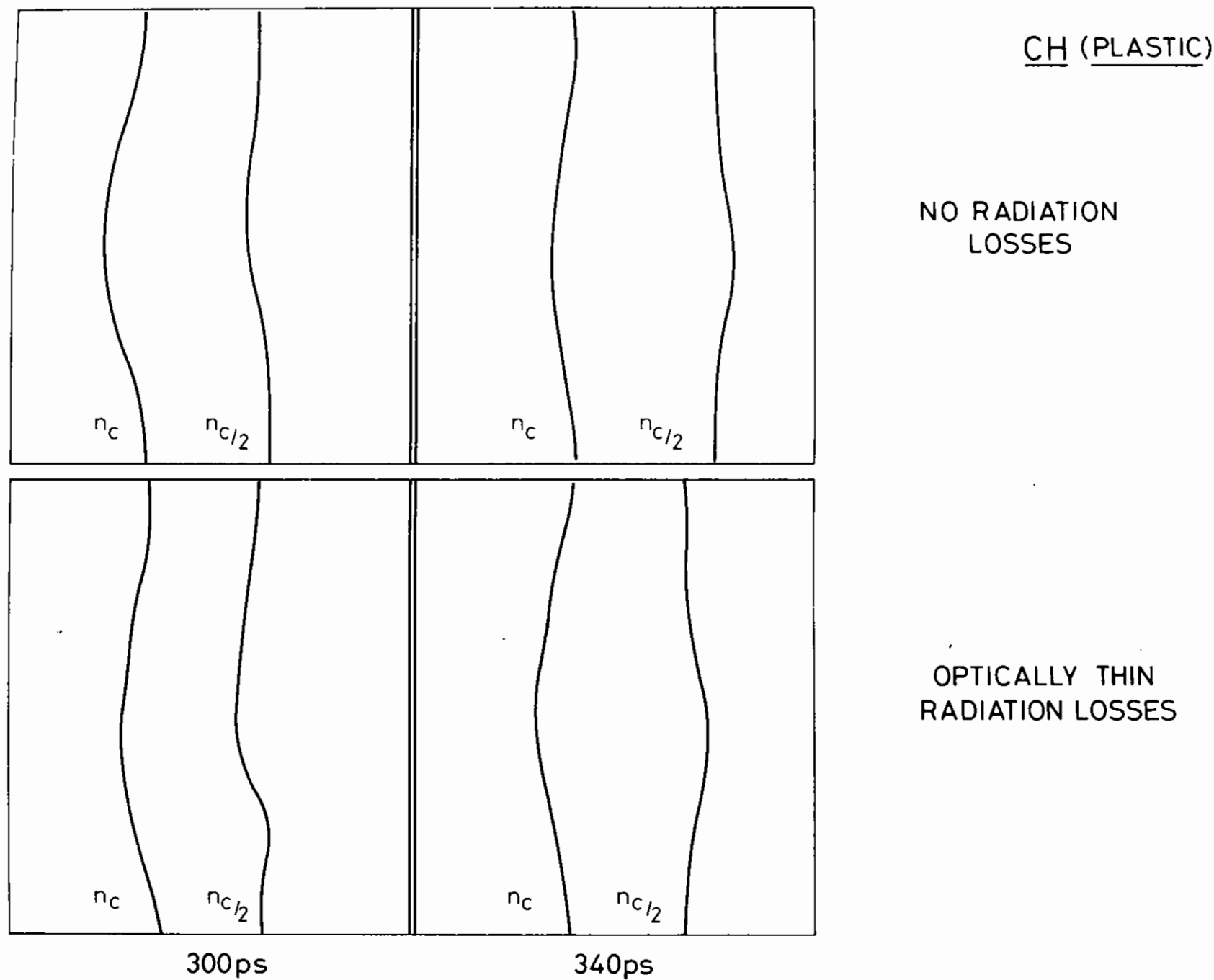


Fig. A5.17 Iso-density contours on a plastic target with and without the radiative loss term.

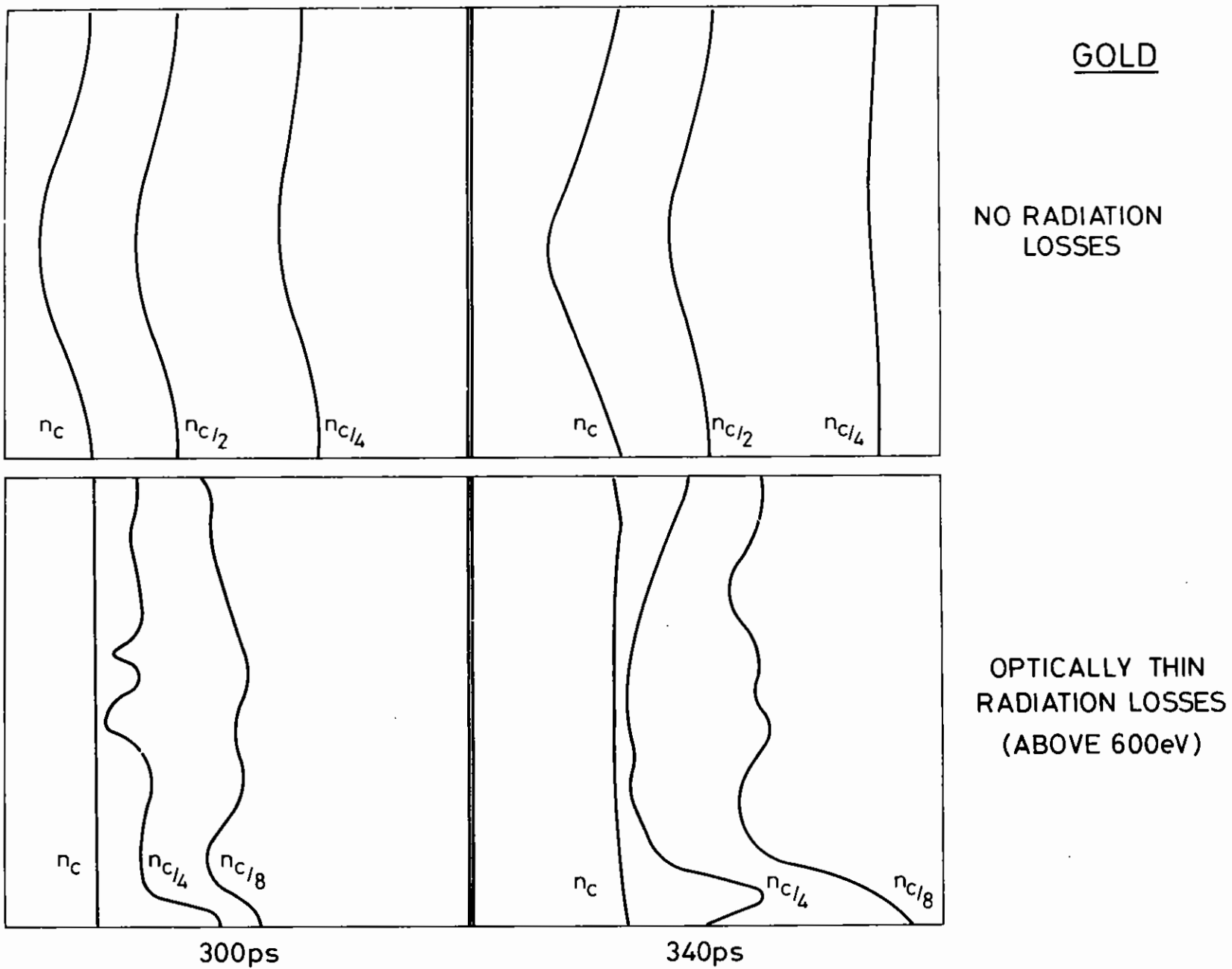


Fig. A5.18 Iso-density contours on a gold target with and without the radiative loss term.

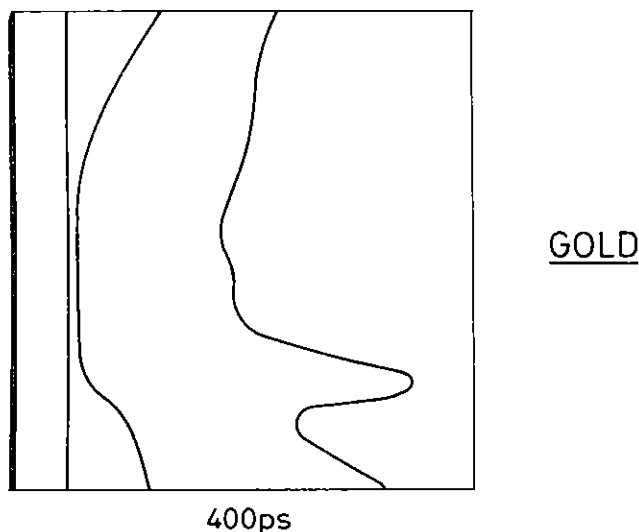


Fig. A5.19 Iso-density contours on a gold target at $t = 400\text{ps}$.

Fig.A5.19 shows the "jet" at 400ps, at which time it has bifurcated but remains confined to the under-dense plasma. The density contours are generally closer together in the simulations including radiation losses since the large radiative losses on Gold targets gives rise to lower temperatures and consequently reduced axial scale lengths.

The distinct lack of symmetry in these results is due to the $(r-z)$ cylindrical mesh having different boundary conditions at $r = 0$ and at $r = r_{\text{max}}$.

Calculations on an intermediate Z material (Silicon) show "jet" structure to a rather lesser extent than on Gold, Fig A5.20, but still noticeable and still confined to the underdense plasma. The temperature change due to the radiative losses on Silicon is clearly much less than for the Gold target so the "jetting" can be attributed directly to the radiative loss instability rather than being a side effect of the reduced temperature. Fig. A5.21 shows calculations on a solid Silicon target showing clearly the same "jet" features even in the absence of target acceleration.

Experimentally the strong plasma "jets" are observed on high Z materials (Gold, Tungsten), and on moderate Z materials (Aluminum). The dimensions and growth times of the jet-like features in the simulations are consistent with the experimental observations.

A5.3.3 Hydro-Code Development

G J Pert (Hull)

The recognition of heat transport from the plasma heated by laser absorption to the ablation surface as a parameter governing the complete hydrodynamic behaviour of the expansion of carbon fibres clearly indicated that a better description of the deposition process was necessary than had been used hitherto. Furthermore such an analysis should allow estimates of the absorptivity of the fibre to be calculated. In view of the small size of the fibre and the large density gradients during the heating phase, it is clear that plasma refraction must play a significant role. On the other

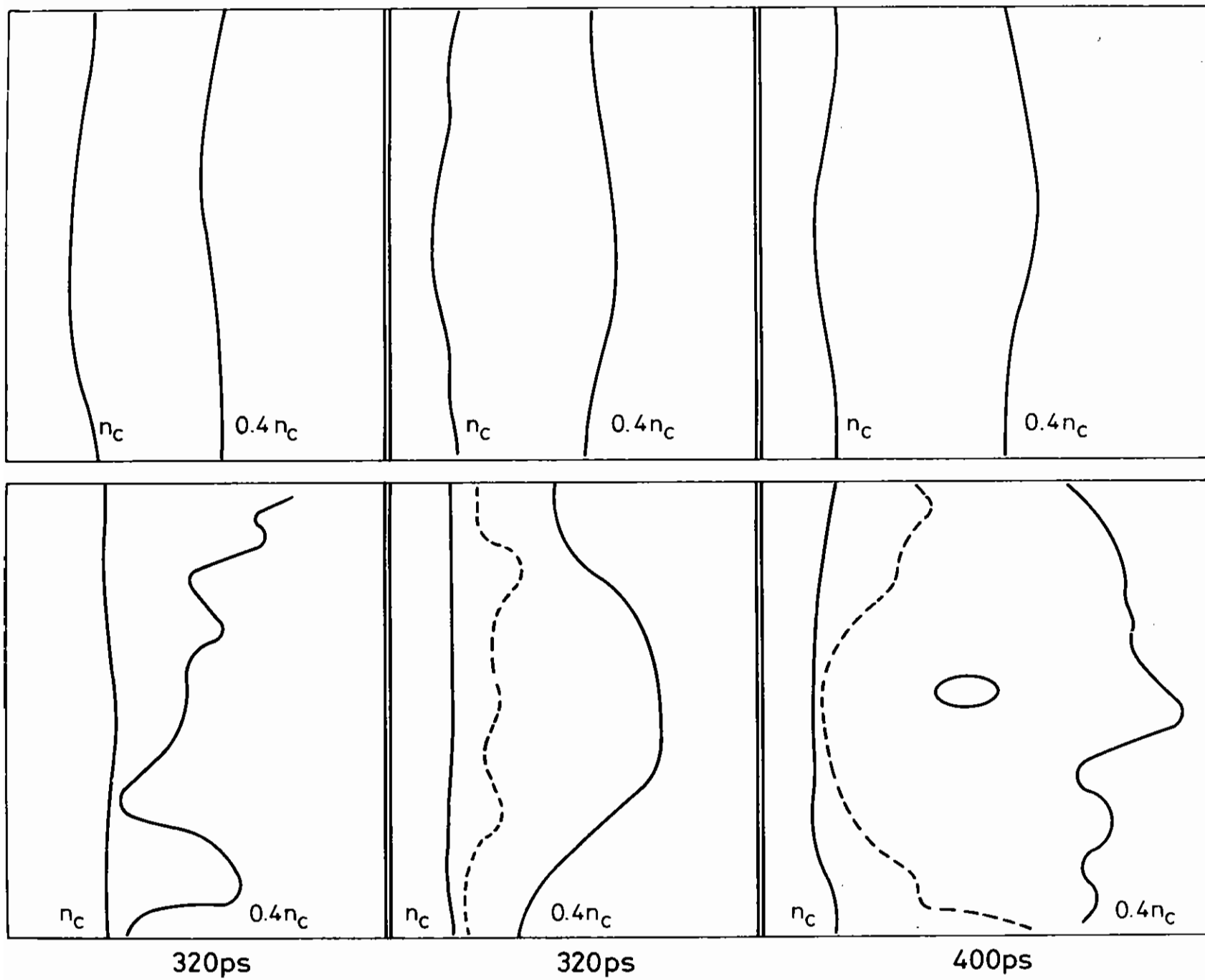


Fig. A5.20 Iso-density contours on a silicon target with and without the radiative loss term.

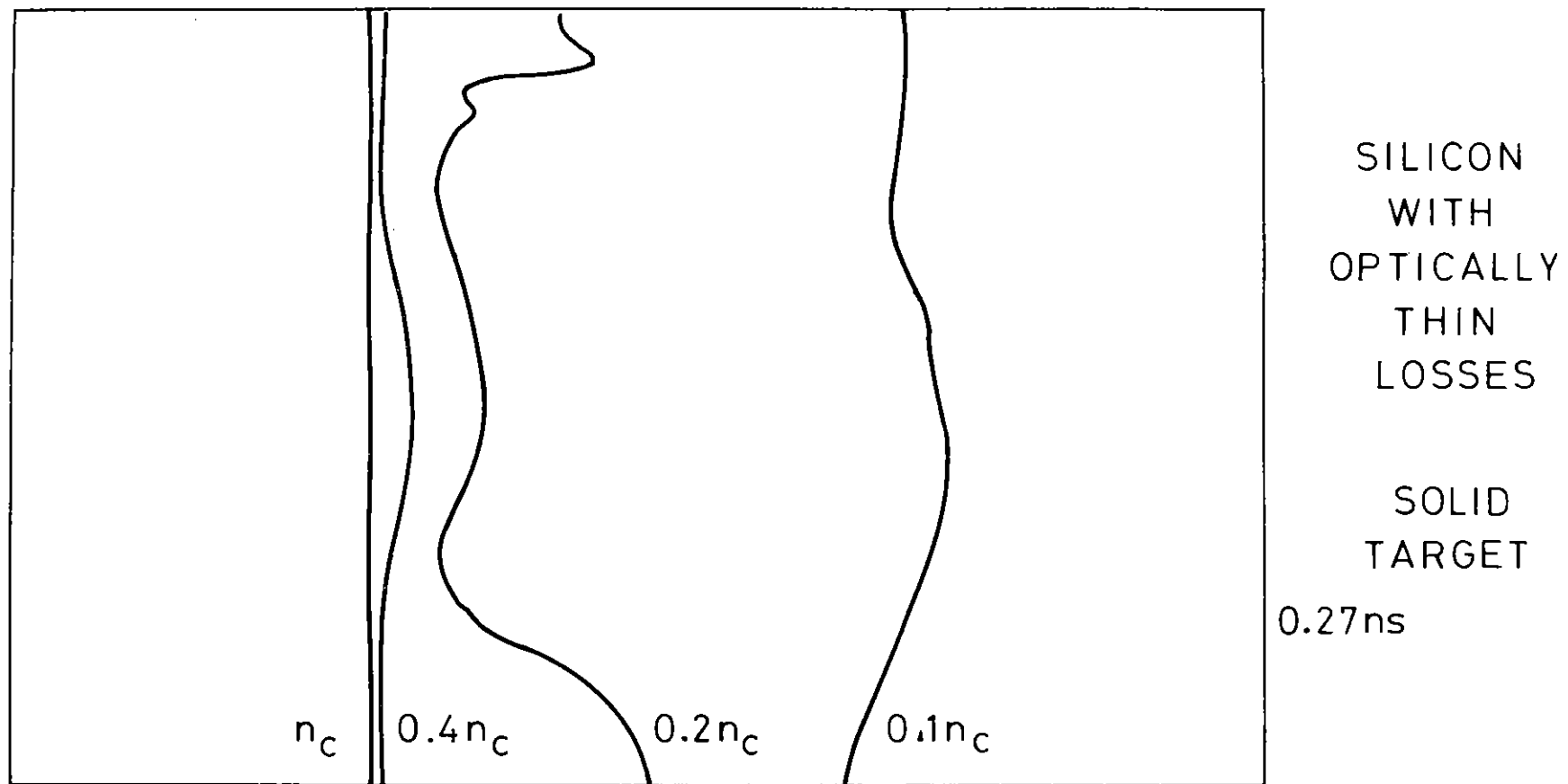


Fig. A5.21 Iso-density contours on a solid silicon target with the radiative loss term.

hand the scale lengths remain large compared to the laser wavelength, so that diffraction and wave effects are unimportant except possibly in the vicinity of the critical surface, where a match to the analytic solution can be used (A5.31). To this end therefore a geometrical optics, ray tracing approach with appropriate account of resonant absorption, is satisfactory, provided the intensity is sufficiently low that non-linear effects such as parametric instability absorption, fast electron generation or profile steepening are ineffective.

We may approach ray tracing in a hydrodynamic cell code via two distinct approximations:

1. The refractive index (density) in the cell is constant - block or zero order approximation. The ray travels in a straight line and is refracted at the discontinuity at the cell boundary.
2. The refractive index (density) varies continuously through the cell, and across the boundary, i.e. everywhere - spline or higher order approximation. The ray travels along a curved path with no local cusps.

Of these the former is generally the easier to programme, but is clearly less accurate. In addition the block approximation must be discarded to deal with resonant absorption. The second method is both more accurate and more physical. In one dimension it can, in principle, be used to any degree of approximation by the appropriate spline. In multidimensions there exists a simple first order spline of wide generality.

In general ray tracing has in the past proved computationally expensive as a large number of rays have had to be followed through the mesh. To alleviate this problem it is important firstly that the ray tracing be carried out in an efficiently written section of code, and that the assignment of absorbed energy be arranged to minimise the required number of rays.

The problems envisaged in this work are hydrodynamically two dimensional although the laser rays themselves may not lie entirely in this plane. The non-planar element can be taken into account by an appropriate modification of the local refractive index as described in ref.A5.31: it will be assumed

that this correction has been made and that the problem reduced to a fully two-dimensional one.

One Dimensional Cylindrical (or Spherical) Systems

In this case the refractive index may be assumed to have either a block or continuous form. Although as already noted the latter allows a more accurate form, the complex nature of the path integrals involving elliptic integrals (A5.31) do not lend themselves to rapid calculation. Despite considerable effort it has had to be conceded that in this case the computational simplicity of the block approximation far outweighs the other advantages.

The block approximation is extremely convenient to apply to calculations in mesh centred Lagrangian code where the refractive index $\mu_{i+\frac{1}{2}}$ is specified at the cell centre, and the boundaries R_i are known. In cylindrical or spherical geometry Snell's law takes the form:

$$\mu \sin \theta = B$$

where θ is the angle the ray makes with the normal, and B is the impact parameter. In passing through the cell from the outer to the inner (say) face the angle changes from $\sin^{-1} B/(\mu_{i+\frac{1}{2}} R_{i+1})$, to $\sin^{-1} B/(\mu_{i-\frac{1}{2}} R_i)$, and on passage across the boundary at R_i to $\sin^{-1} B/(\mu_{i-\frac{1}{2}} R_i)$. If $B > \mu_{i-\frac{1}{2}} R_i$ the ray is reflected at that boundary and re-enters cell i at angle $\sin^{-1} B/(\mu_{i+\frac{1}{2}} R_i)$.

This completes the ray trace.

The path length across the cell is easily calculated, since the path is a straight line, to be

$$Z = \sqrt{R_i^2 + R_{i+1}^2 - 2 \{B^2 + [R_i^2 - B^2/\mu_{i+\frac{1}{2}}^2] \{R_{i+1}^2 - B^2/\mu_{i+\frac{1}{2}}^2\}\}}$$

which may be used to find the inverse bremsstrahlung absorption in the cell.

If reflection occurs at the boundary from the cell i to $(i-1)$ resonant absorption will occur. To account for this the refractive index must be allowed to vary from $\mu_{i-\frac{1}{2}}$ to $\mu_{i+\frac{1}{2}}$ in some continuous form e.g.

$$\mu^2 = 1 - a/r - b/r^2$$

where

$$a = R_{i+\frac{1}{2}} R_{i-\frac{1}{2}} - (\mu_{i+\frac{1}{2}}^2 R_{i+\frac{1}{2}}^2 - \mu_{i-\frac{1}{2}}^2 R_{i-\frac{1}{2}}^2) / (R_{i+\frac{1}{2}} - R_{i-\frac{1}{2}})$$

$$b = -R_{i+\frac{1}{2}} R_{i-\frac{1}{2}} \{1 - (\mu_{i+\frac{1}{2}}^2 R_{i+\frac{1}{2}} - \mu_{i-\frac{1}{2}}^2 R_{i-\frac{1}{2}}) / (R_{i+\frac{1}{2}} - R_{i-\frac{1}{2}})\}$$

The turning point is given by

$$R = \frac{1}{2} \{ (R_{i+\frac{1}{2}} + R_{i-\frac{1}{2}}) - (\mu_{i+\frac{1}{2}}^2 R_{i+\frac{1}{2}}^2 - \mu_{i-\frac{1}{2}}^2 R_{i-\frac{1}{2}}^2) / (R_{i+\frac{1}{2}} - R_{i-\frac{1}{2}}) \}$$

$$+ \sqrt{\{ (R_{i+\frac{1}{2}} - R_{i-\frac{1}{2}}) - (\mu_{i+\frac{1}{2}}^2 R_{i+\frac{1}{2}}^2 + \mu_{i-\frac{1}{2}}^2 R_{i-\frac{1}{2}}^2) / (R_{i+\frac{1}{2}} - R_{i-\frac{1}{2}}) \}}$$

$$+ 4 \{ B^2 - \mu_{i+\frac{1}{2}}^2 \mu_{i-\frac{1}{2}}^2 R_{i+\frac{1}{2}}^2 R_{i-\frac{1}{2}}^2 / (R_{i+\frac{1}{2}} - R_{i-\frac{1}{2}})^2 \}$$

the root lying between $R_{i-\frac{1}{2}}$ and $R_{i+\frac{1}{2}}$ being chosen. The intensity loss of the ray by resonant absorption is then directly calculated from the refractive index at R by the method indicated in ref. A5.31.

To complete the algorithm a bundle of rays of width appropriate to the focal spot or the plasma are assigned energies according to the intensity distribution in the focal spot and traced through the plasma. The energy loss of each ray is directly assigned to the appropriate cell as it traverses its path. This algorithm is incorporated in the one dimensional Lagrangian code LAGR.

Two Dimensional Cartesian Systems

In this case it is appropriate to use a simple spline to describe the refractive index continuously throughout the medium. The refractive index is known at the mesh points whose co-ordinates themselves are known. In contrast to the usual quadrilateral zone the mesh may be divided into two triangles, $((i,j), (i+1,j), (i+1, j+1))$ and $((i,j), (i,j+1), (i+1, j+1))$ each quadrilateral cell being formed of an upper and lower triangle. If the electron density is assumed to be linear over each triangle with the mesh values at each apex then it is easily seen that the value of the electron density and therefore the refractive index on each boundary depends only on its end points, and is therefore continuous over the mesh.

It is easily shown A5.32 that the ray path in a linear electron density gradient is a parabola, whose intersections with the triangular sub-cell will give the entry and exit points of the ray. The ray path is easily traced by following the ray through each exit point from triangle to triangle. The absorption of energy by the plasma can be assigned between exit points by calculating that due to inverse bremsstrahlung and resonant absorption on each path segment.

Thus far the absorbed energy is identified with the rays only, and must now be assigned to the plasma on the mesh. It is at this stage that care is required to ensure that the algorithm is efficient and does not require an excessive number of rays for accuracy. Since the scheme is continuous it may be plausibly assumed that any ray lying between two marked rays will in a general sense have a trajectory which conforms to the two marked rays. Thus we may regard the calculated ray paths as boundaries of a tube of energy flow, the energy within the tube being absorbed according to a mean of the absorption along the bounding rays. The assignment of absorbed energy to the cell is now straightforward, being that fraction of the tube energy which lies within the cell boundary and is absorbed within the cell. Although in principle simple, the actual detailed prescription is somewhat complicated and is given in ref. A5.32. In this way relatively few rays can be used, without significant error.

This algorithm for ray tracing and energy deposition was developed some years ago A5.32 as a post-processor code. The programmatic form of this code was not suitable for use as an internal sub-routine being far too slow. The algorithm was recently re-written into a more efficient code and included in versions of POLUX. The overheads in using ray tracing in the code are relatively small, both in time (~ 20% increase) and store. Full details of the numerical form will be given in a future publication.

5.3.4 Effects of Beam Spot Size in the performance of HIF targets

T D Beynon and E H Smith (Birmingham)

If one postulates a uniform irradiance of an ICF target via a uniform distribution of beams, then using a 1D hydrocode, such as MEDUSA, it is possible to examine the effects of a) the width of each beam and b) the radial intensity distribution in each beam. Such studies have been performed for ion-beam driven targets with MEDUSA using gaussian-distributed beam densities of the form $\exp(-x^2/2\sigma^2)$ where x is measured from the centre of a cylindrical beam orthogonal to the direction of propagation. The spot size of a beam is defined as σ .

A heavy ion beam driven target is considered with a full dE/dx treatment. The heavy ion beams are assumed to consist of uranium ions of such energy as to just be able to penetrate initially to within 1.1mm of the target's centre. The energy of the ions is then kept constant and the range is therefore allowed to vary. The pulse shape of the ion beam consists of a pre-pulse followed by a main pulse. Both are of constant power, the main pulse having either a power level of $2.5 \times 10^{14}W$ or $1.0 \times 10^{15}W$. The pre-pulse is set to 0.5% of the main power pulse and lasts until either the shock wave has reached to within 0.85mm of the target's centre or 20ns has elapsed.

Five different spherical targets have been used in this study and are described in Fig.A5.22. All have identical central structures, a 0.8mm D-T gas filled hole surrounded by 0.2mm of solid D-T which is in turn surrounded by a 0.1mm ($1.13kgm^{-2}$) lead shell. The thickness of the lead shell was chosen to give it an aspect ratio of 10 which is considered to be consistent with Rayleigh Taylor stability. This standard core is surrounded by either simple shells, $2.26kgm^{-2}$ of lead or carbon, or $1.13kgm^{-2}$ of lithium deuteride (LiD), or two-layer shells of either lead and carbon or carbon and LiD. The LiD was made thinner than the carbon in order to maintain a reasonable total target size. The targets chosen are not intended to represent optimized ICF targets but rather to illustrate the beam-target coupling effects.

A5.32

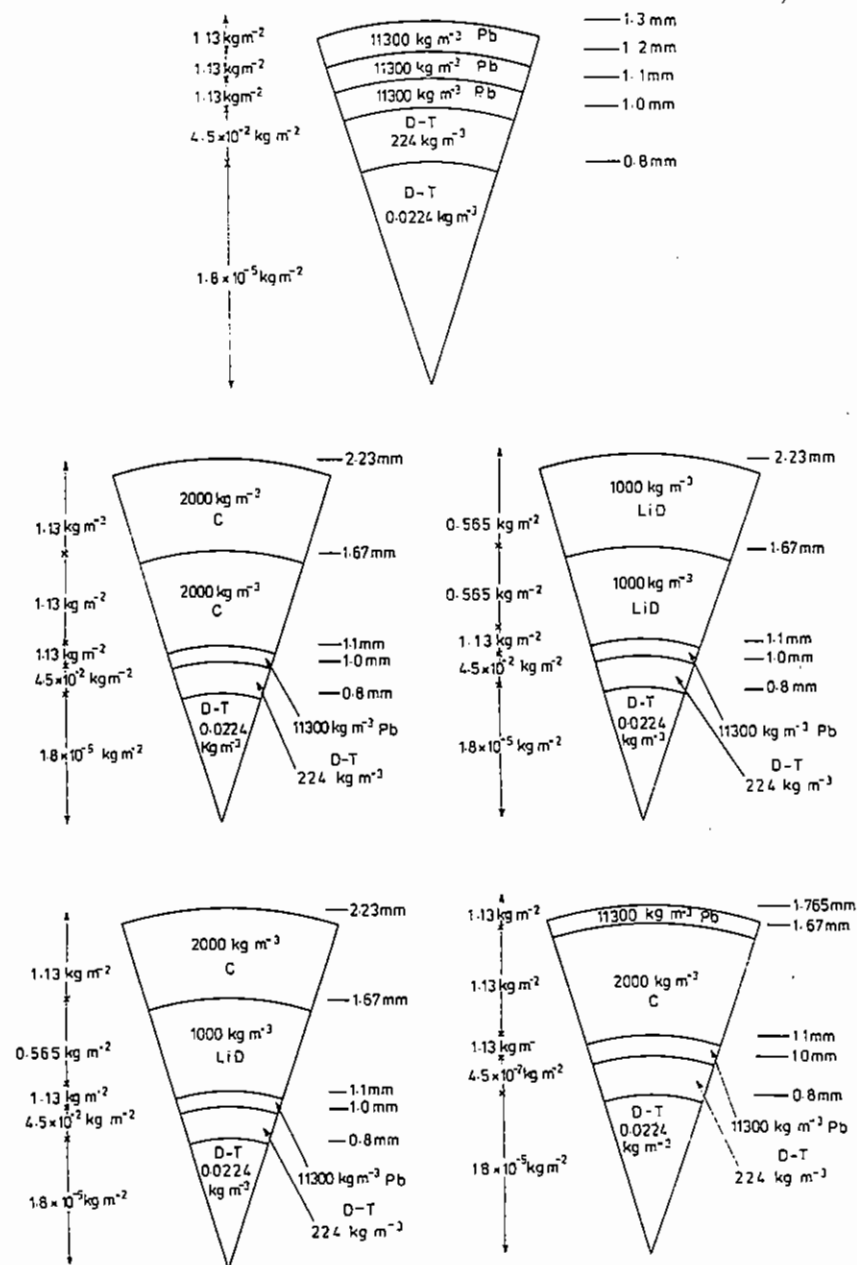


Fig. A5.22 Various target configurations used in the present study.

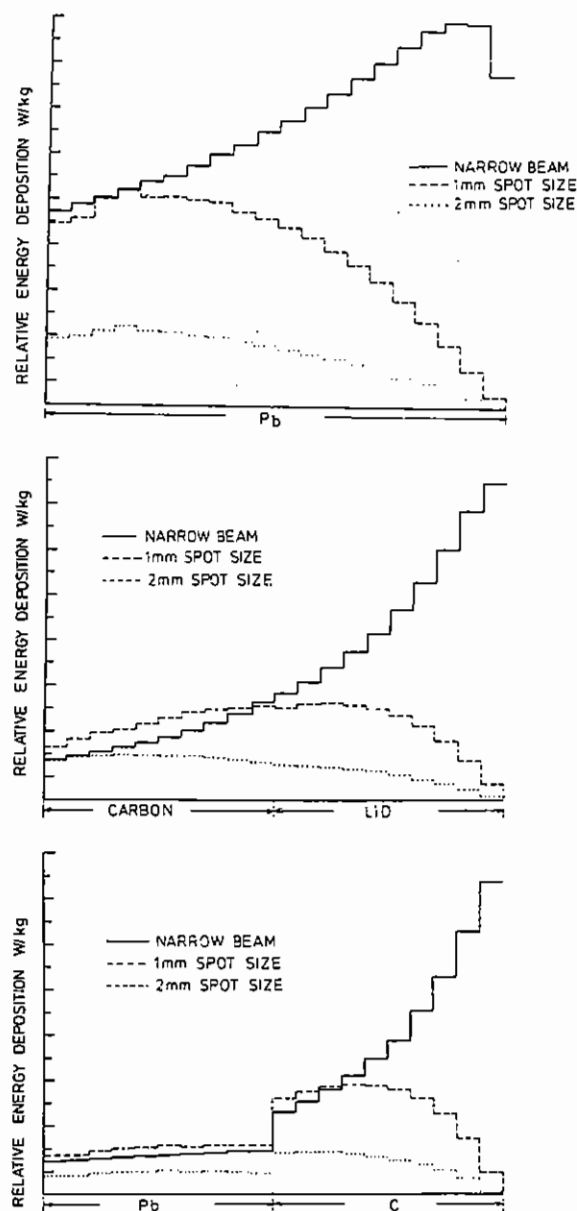


Fig. A5.23 Energy deposition profiles for targets (a), (d) and (e) (cf. Fig. A5.22). In each case the beam is travelling from left to right.

In Fig.A5.23 deposition profiles are shown for a selection of the targets for different spot sizes. As can be seen the energy deposition rates for narrow beams ($\sigma \sim 0$) and those of size ($\sigma \sim 1\text{mm}$) are similar near the target surface. The 1mm beams even tend to give slightly higher depositions near the target. For the 2mm spot size beams however the effect of the beam missing the target predominates and the deposition is noticeably reduced even near the surface. Further into the target the energy deposition profiles of the narrow beams rise significantly to peaks near the end of their ranges. The beams of finite spot size however give profiles which decline appreciably towards the end of their ranges.

In this case only those ions at or near the centre of each beam are able to penetrate far into the target. During the early portion of the pre-pulse the range of the ions shortens, especially in the denser shells, due to both the temperature and density effects. The time-averaged deposition profile due to the pre-pulse will therefore show a region which has been heated by the ions for a short time only. For a narrow beam this region will be heated considerably more than for a wider beam.

The coupling efficiency parameter, the 'maximum inward kinetic energy' (MKE), is defined by

$$\text{MKE} = \max_{\text{time}} \ln \left\{ \sum_{i=1,n} \frac{1}{2} m_i v_i^2 H(-v_i) \right\} \quad (1)$$

where m_i and v_i are the masses and velocities of the n numerical cells used in the simulation code MEDUSA, inward velocities being negative. H is the Heaviside step function. The maximum value is achieved when the lead shell starts to slow down due to the main reflected shock reaching its inner surface. MKE is typically only 1 or 2% of the total energy delivered in the beams. This low value is due to the main pulse being left on till the end of the simulation so that only part of the energy delivered is utilized. The conservative aspect ratio of the lead shell also contributes to this low efficiency.

The simplest target is that described in Fig.A5.22(a). The MKE for this

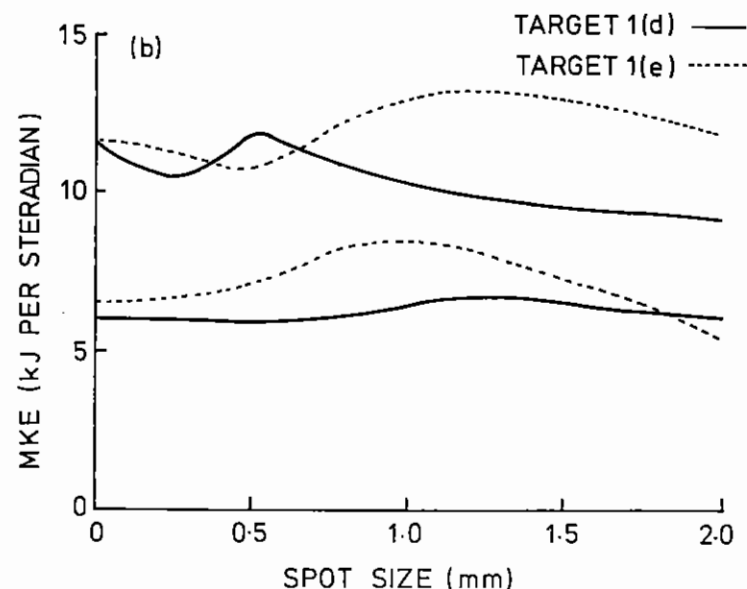
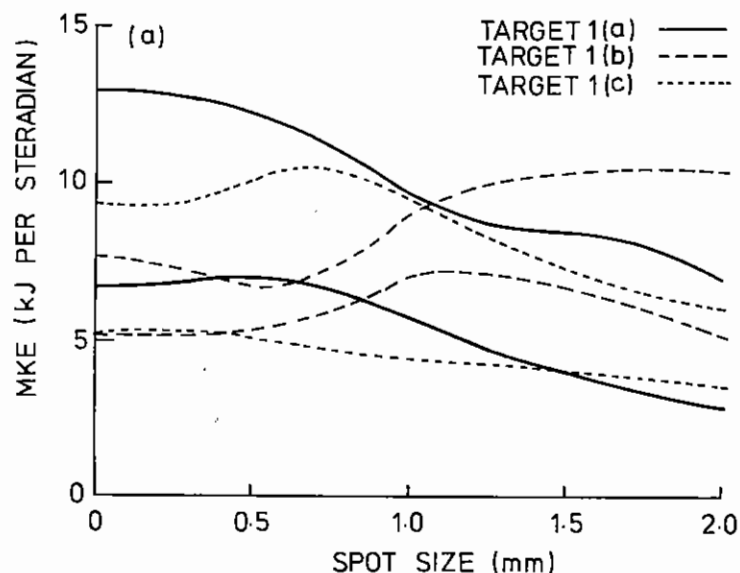


Fig. A5.24 The MKE in the targets of Fig. A5.22 for power levels of 2.5×10^4 W (lower case for each pair) and 1.0×10^5 W (upper case for each pair).

target is shown in Fig.A5.24(a). As would be expected the MKE drops with increasing spot size. This can be attributed to both the larger fraction of the beam missing the target at the larger spot sizes and the absorbed energy being predominantly deposited near the surface where it produces a very poor driving force for the target shell (cf Fig.A5.23(a)). However at the higher power level, at spot sizes of about 1.5mm, the fall in MKE becomes much more gradual producing a hump in the curve. This arises since a cool region is left behind by this wide beam when the range shortens. This cool region is now appreciably compressed by the main pulse. Because the thermal conduction is high enough this region is warmed during the latter part of the main pulse, due both to the conduction and the beam burn-through it produces, giving high pressures and thus higher values of MKE.

The other two simpler targets described in Figs.A5.22(b) and A5.22(c) show these effects more strongly as can be seen in Fig.A5.24(a). In particular, for the carbon target the inward heat transfer can be very large during the main pulse. The LiD target shows less of these effects, partially due to less range shortening and also due to very large heat conduction within the lithium with little coupling into the lead. This both reduces compression of the inner layers of the lithium and the heating of the lead so that the production of a hot dense high pressure region is reduced.

For very small spot size beams the ion beam is able, in the higher-power carbon shell case, to burn an appreciable distance into the lead because of the high temperature resulting from the high deposition at the end of the beam range. At later times, when the burn-through has become appreciable, the target is becoming well compressed and the beam can reach to within 0.5mm of the target centre in the middle of the lead shell. This produces a great deal of radial focussing which produces high temperatures and pressures, giving a rise in MKE.

The two more complex targets described in Figs.A5.22(d) and A5.22(e) have MKE profiles, Fig.A5.24(b), similar to those produced by shells composed solely of their inner layers. The MKE tends however to be slightly larger due to the tamping effects of the outer denser layers. For the carbon-LiD target the higher compressions achieved due to the better coupling enable radial

focussing and burn-through to occur at small beam widths. The lead-carbon shell shows somewhat faster deterioration at large spot sizes compared with the pure carbon shell. This is due to the large fraction of the beam energy which is deposited on the outside of the lead shell not being able to conduct through the lead and reach the inner carbon regions of the shell.

We see that beam spot sizes do not effect target performance in any simple manner. The effects depend on the power, the particular target configuration and the effects of the beam range changes. As a result no optimised target can be designed without knowledge of both the energy and the quality of the beam. If these parameters are not known the choice of a shell structure to give an optimised value of MKE is impossible.

Another important feature of these results is that MKE does not scale with power. For a factor 4 increase of power MKE generally rises by less than a factor of 2. If, for a given target, larger values of MKE are required for ignition, due either to less than perfect compression or ignition difficulties, dramatic increases in beam powers might be needed.

A5.4 Radiation Transport and Atomic Physics

A5.4.1 A Ray Tracing Code for Radiation Transport in MEDUSA

A C White (Glasgow)

A program has been written which obtains solutions of the equation of radiative transfer (ERT)

$$dI/ds = k (B - I) \quad (1)$$

in the case of spherical symmetry. In the above, $I = I(r, t, \nu, \Omega)$ is the specific intensity (Joules/sec/metre²/Hz/ster). B is the source function, not necessarily Planckian but assumed given. k is the opacity (m⁻¹), also assumed given. ds is the incremental distance along the ray path. As is well known, the deceptively innocuous equation (1) describes a 'heavy' problem, as revealed by the argument list of the unknown quantity, I . Ideally we should solve for the time-dependent specific intensity at all

radial points, r , at all frequencies, ν , and in all directions, Ω . In this note, we shall neglect the frequency dependence (loosely termed the 'grey' problem), though a multi-group treatment will finally appear. We are also tackling the time-independent problem, insofar as there is no time derivative in (1). We shall assume that within any hydrodynamic timestep, the local plasma conditions have adjusted very rapidly. This common assumption implies that the radiation has only a 'first-order' influence on the prevailing plasma state, and effectively linearizes the problem. It is, of course, sometimes open to question, whenever radiation forms a significant part of the energy transfer.

With these assumptions and in the context of the MEDUSA system, the problem reduces to solving the ERT across concentric shells that are each spatially uniform (Fig.A5.25).

Energy balance: To calculate the radiative contribution, we shall be concerned with angle integrated quantities at each of the boundaries, viz. radiative energy density, and flux, given respectively by:

$$E = 1/c \int_{4\pi} I d\Omega \quad F = \int_{4\pi} \Omega I d\Omega$$

Stated otherwise, we wish to relate the energy, flux, etc. at any given boundary to that at adjacent boundaries. The problem is of course incomplete without specifying the flux incident at the outermost boundary. We assume no external flux is incident. The radiation-electron energy exchange within any shell comes simply from the difference in fluxes into and out of the shell. The balance of radiative energy is added to or subtracted from the plasma electrons. To complete our description of the coordinate system (Fig.A5.26) we note that the specific intensity must be independent of the azimuthal angle ϕ . Angle θ , however, between the radius vector and Ω , will still enter the treatment.

Method of solution: To solve for $I(r, \theta)$ at any given boundary, there are essentially three situations to be considered (see Fig.A5.27). What all of these have in common, is that their formal solution is in each case given by

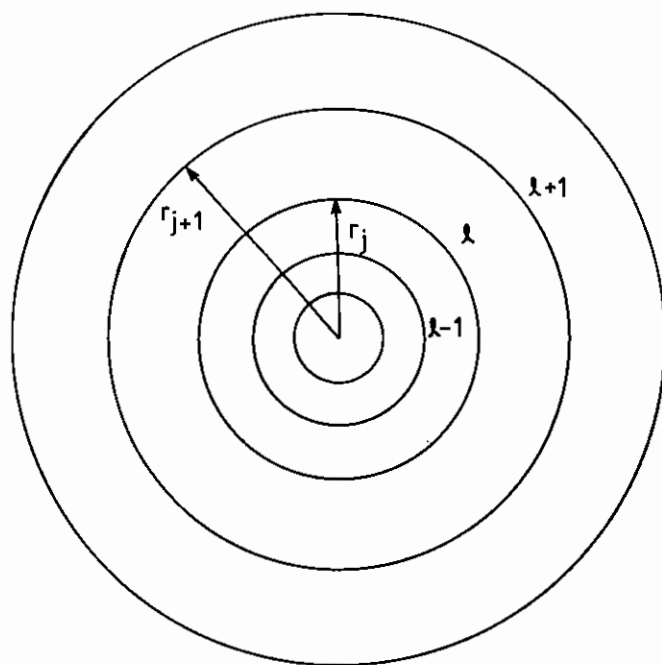


Fig. A5.25 Radiation transport across concentric shells. Opacity and source functions are assumed specific for each of the shells.

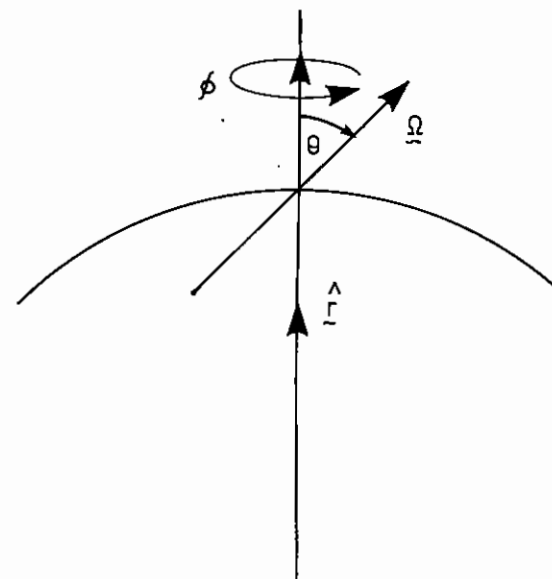


Fig. A5.26 Co-ordinates at a typical boundary. Spherical symmetry precludes any ϕ dependence.

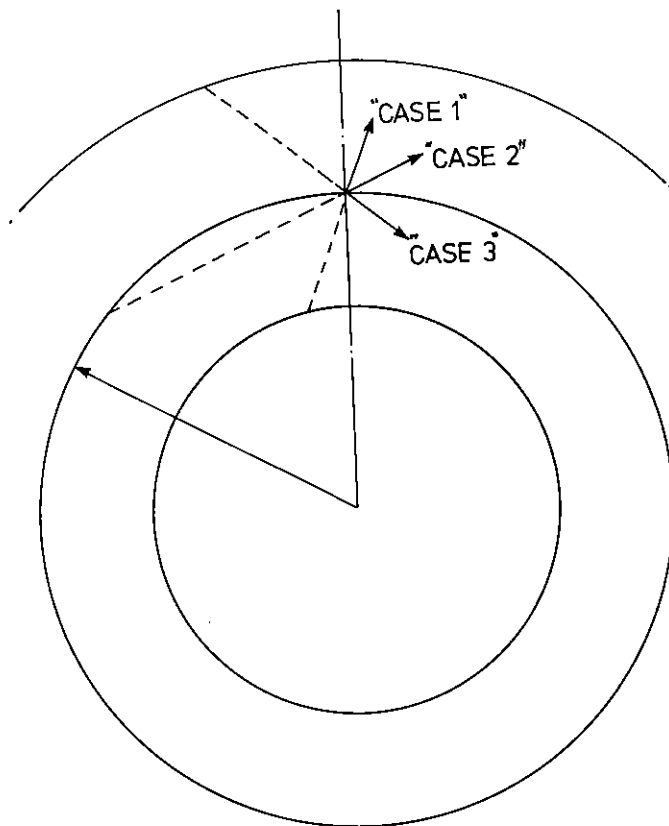


Fig. A5.27 Rays to be considered at a typical boundary.
In 'case 1', rays emerge from an interior boundary.
In 'case 2', rays impinge from an exterior boundary.

$$I(P, \theta) = I(Q, \psi) \exp(-kx) + B (1 - \exp(-kx)) \quad (2)$$

- the integration of equation (1). This latter equation describes the specific intensity at point P, at angle θ in terms of (i) the attenuated flux that left point Q, at angle ψ to the radius vector, and (ii) the energy added to the ray by the emitting matter between points P and Q (distance x). Because the angular variable θ will be discretised, and interpolation will have to be used to obtain the intensity at angle ψ , it is a feature of this model that we can solve for "case (3)" (see figure A5.27) angles straightaway, working inwards from the outermost boundary. "Case (2)" angles are readily derived from these results. Finally we solve for "Case (1)" angles from the centre outwards. Despite the coupling of intensities at adjacent boundaries, no iteration is required, even for interpolated values.

Choice of angles: It should be apparent that one of the strengths of this approach is that equal importance is assumed for each boundary, i.e. the specific intensity is calculated at chosen values of θ , the so-called nodal angles irrespective of radial position. This distinguishes it from other ray tracing methods (Ref. A5.33, A5.34). Its weakness stems from the interpolation mentioned above. There are two guides to the choice of nodal angles: (i) their number. If there are many, the interpolation is improved, at the expense of computer time. (ii) their position. Since we integrate over this angular variable, we might choose nodal angles that optimise a quadrature scheme. Currently, the code uses equally spaced nodal angles, however.

Preliminary Test Results: The program was tested for the case of a uniform sphere, i.e. constant absorption coefficient and source function. The results in this case are known analytically, as a function of r and θ . The numerical results reproduce the analytical results to within about 1% when 20 (equally spaced) angles are used, and to within about 5 or 6% when 10 angles are used. The nearer θ is to 0 or π , then the more exact are the results (which is not too surprising since at these extremities the exact solution is obtained).

A more extensive report of this code is in preparation, together with some

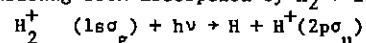
results of its implementation in MEDUSA. Work is being carried out to develop an efficient time-dependent ionisation package which, inter alia, should provide better estimates of the local opacities and source functions mentioned above.

A5.4.2 Two-centre electronic wave functions in plasmas

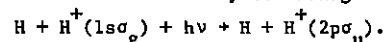
S J Rose (RAL)

One of two general approaches is usually used in the calculation of the electronic structure of a plasma. The first involves an initial calculation of the electronic structure of an isolated ion. The effect of the surrounding plasma is introduced by perturbation theory. In the second method the effect of the surrounding plasma is included in an initial calculation of a spherically symmetric (such as Thomas-Fermi) potential. The electronic structure is then calculated by solving the electronic wave equation using this potential. Such single-centre calculations are regularly performed to obtain for example, the radiative opacity of a plasma.

Bates (A5.36) has investigated the contribution to the radiative opacity of stellar atmospheres arising from absorption by H_2^+ ions.

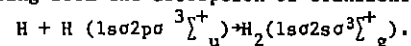


and from the transient molecule formed by colliding H atoms and H^+ ions.



Although Bates (A5.36) calculates the contribution to the stellar opacity from these processes, because of lack of data on the H_2^+ ion, in each case he only considers the transition $1s\sigma_g \rightarrow 2p\sigma_u$.

Several authors (A5.38, A5.43, A5.44) have considered the contribution to the stellar opacity arising from the absorption of transient molecular hydrogen



However they only consider a bound final molecular state.

Subsequent applications of the concept of transient molecule formation in a plasma have been concerned with the calculation of line shapes and in particular quasi-molecular satellites (e.g. A5.42)

It is the purpose of the present note to point out that the use of two-centre wave functions to describe the electronic structure of a plasma predicts contributions to plasma properties not existing in the usual single-centre treatment. These contributions may be described as arising from the 'transient chemistry' of the plasma.

Transient H_2^+ formation in a hydrogen plasma

As in the work of Bates (A5.36) this note assumes the validity of the Born-Oppenheimer approximation which allows the solution of the electronic wave equation at fixed nuclear separations; the nuclear equation of motion involving the potential energy surface obtained. To be of most use in this application, the two-centre electronic wavefunctions must be known accurately. The molecule for which the most accurate and extensive electronic calculations have been performed is the hydrogen-molecule ion, H_2^+ . The Schrodinger equation for this system has been solved exactly (in the Born-Oppenheimer approximation) for several bound states (A5.37) and the transition energy and dipole matrix element have been calculated for a range of internuclear separations for several transitions (A5.39, A5.41). In the present note, two H_2^+ electronic transitions, $1s\sigma_g \rightarrow 2p\sigma_u$ and $1s\sigma_g \rightarrow 2p\pi_u$, are studied which exemplify features of transitions between two-centre electronic wavefunctions. The transient H_2^+ ions may be considered to form in binary collisions between H atoms and H^+ ions in a (low temperature) hydrogen plasma. We shall not use the treatment of Bates (A5.36) to determine an absorption opacity, but in order to examine the physical processes occurring, we shall consider the molecular oscillator strength $f_{i \rightarrow f}$. This is given (in atomic units) by

$$f_{i \rightarrow f} = \frac{2}{3} (\epsilon_f - \epsilon_i) g_{i \rightarrow f} |I_{i \rightarrow f}|^2$$

where $\epsilon_f - \epsilon_i$ is the energy difference between initial and final states, $g_{i \rightarrow f}$ is the degeneracy factor (A5.40) and $I_{i \rightarrow f}$ is the dipole matrix element

for the transition

$$I_{i \rightarrow f} = \langle \psi_i(\underline{r}) | \underline{r} | \psi_f(\underline{r}) \rangle$$

with the initial and final electronic wavefunctions given by $\psi_i(\underline{r})$ and $\psi_f(\underline{r})$

respectively. The values of $\epsilon_f - \epsilon_i$, $I_{i \rightarrow f}^2$ and $f_{i \rightarrow f}$ are tabulated for a range of internuclear separations in tables A5.1 and A5.2 and the oscillator strengths are plotted in figure A5.28 to show their dependence on internuclear separation.

For the $1s\sigma_g \rightarrow 2p\pi_u$ transition, as the nuclei separate the oscillator strength decreases to zero. This is because in the separated atom limit this corresponds to the atomic 'transition' $1s \rightarrow 1s$. The influence of the plasma is therefore to introduce an absorption mechanism not seen in the isolated atom and which could not have been predicted by first-order perturbation of the hydrogen atom by the neighbouring proton. It is necessary to consider the solution of the Schrodinger equation in a two-centre system to obtain such 'bonding' and 'antibonding' solutions. Transition of an electron between 'bonding' and 'antibonding' orbitals by absorption of a photon rearranges the electronic distribution between the nuclei. Because, in the limit of large internuclear separation the transition energy tends to zero, this results when the nuclei separate, in the transfer of that energy to the kinetic motion of the ions. This coupling between radiation and ionic kinetic energy is evident only the solution of the electronic wave equation in a two-centre system.

The molecular transition $1s\sigma_g \rightarrow 2p\pi_u$ corresponds to the atomic transition $1s \rightarrow 2p_{\pm 1}$ in both united and separated atom limits and the oscillator strength is in each case 0.278. The ratio

$$(f_{1s\sigma_g \rightarrow 2p\pi_u} - f_{1s \rightarrow 2p_{\pm 1}}) / f_{1s \rightarrow 2p_{\pm 1}}$$

is, therefore, a measure of the contribution to the oscillator strength from molecular effects and is tabulated in table A5.1. It can be seen to be significant even for large values of the internuclear separation. Although in the separated atom limit, the energy difference between $1s\sigma_g$ and $2p\pi_u$ is

Table A5.1 $1s\sigma_g \rightarrow 2p\pi_u$ H_2^+ oscillator strengths

Internuclear separation R^*	$ I_{1s\sigma_g \rightarrow 2p\pi_u} ^2$	$\epsilon_{2p\pi_u} - \epsilon_{1s\sigma_g}^{*†}$	$f_{1s\sigma_g \rightarrow 2p\pi_u}$	$\frac{f_{1s\sigma_g \rightarrow 2p\pi_u} - f_{1s \rightarrow 2p_{\pm 1}}}{f_{1s \rightarrow 2p_{\pm 1}}}$
0 ⁺	0.139	1.500	0.278	0.0
1	0.302	0.978	0.393	0.42
2	0.512	0.674	0.460	0.66
3	0.684	0.524	0.479	0.73
4	0.784	0.445	0.466	0.63
5	0.800	0.403	0.430	0.55
6	0.755	0.382	0.384	0.38
7	0.685	0.372	0.339	0.22
8	0.619	0.368	0.304	0.09
9	0.568	0.367	0.278	0.0
10	0.534	0.368	0.262	-0.06
11	0.514	0.369	0.253	-0.09
12	0.505	0.370	0.249	-0.10
13	0.504	0.371	0.249	-0.10
14	0.508	0.372	0.252	-0.09
15	0.514	0.373	0.256	-0.08
16	0.521	0.374	0.260	-0.06
17	0.528	0.374	0.264	-0.05
18	0.534	0.375	0.267	-0.04
20	0.544	0.375	0.272	-0.02
25	0.553	0.375	0.277	0.0
30	0.555	0.375	0.277	0.0
∞^{++}	0.555	0.375	0.278	0.0

* In atomic units

+ H_2^+ transition $1s \rightarrow 2p_{\pm 1}$

++ H transition $1s \rightarrow 2p_{\pm 1}$

† Values taken from Madsen and Peek (A5.39) (except $R = 0$ and ∞)

‡ Values taken from Ramaker and Peek (A5.41) (except $R = 0$ and ∞)

Table A5.2 $1s\sigma_g \rightarrow 2p\sigma_u$ H_2^+ oscillator strengths

Internuclear separation R^*	$ I_{1s\sigma_g \rightarrow 2p\sigma_u} ^{2*†}$	$\epsilon_{2p\sigma_u} - \epsilon_{1s\sigma_g}^{*†}$	$f_{1s\sigma_g \rightarrow 2p\sigma_u}^{\ddagger}$
0^+	1.39×10^{-1}	1.50	1.39×10^{-1}
1	4.55×10^{-1}	8.87×10^{-1}	2.69×10^{-1}
2	1.10	4.35×10^{-1}	3.20×10^{-1}
3	2.05	2.09×10^{-1}	2.87×10^{-1}
4	3.50	1.01×10^{-1}	2.35×10^{-1}
5	5.57	4.71×10^{-2}	1.75×10^{-1}
6	8.29	2.13×10^{-2}	1.18×10^{-1}
8	1.54×10^1	3.96×10^{-2}	4.07×10^{-2}
10	2.45×10^1	6.78×10^{-4}	1.11×10^{-2}
15	5.60×10^1	6.88×10^{-6}	2.56×10^{-4}
20	9.98×10^1	5.17×10^{-8}	4.10×10^{-6}
∞^{++}	0.0	0.0	0.0

* In atomic units

+ He^+ transition $1s \rightarrow 2p_0$

++ H 'transition' $1s \rightarrow 1s$

† Values taken from Madsen and Peek (A5.39) (except $R = 0$ and ∞)

‡ Values taken from Ramaker and Peek (5.41) (except $R = 0$ and ∞)

§ Some of the oscillator strengths differ from the (less accurate) calculations of Bates (A5.35)

0.375 a.u., there exist internuclear separations (table A5.1) for which the energy difference is less than this value. For photon absorption at such a separation in the transient molecule, the excitation energy as the nuclei separate is partly made up from the absorbed photon energy and partly from some fraction of the initial kinetic energy of the ions. If the photon energy is greater than the excitation energy as the nuclei separate, then the excess energy appears as ionic kinetic energy.

At high densities, the binary collision approximation fails. Treating the plasma by a similar technique to that described above would require multi-centre wavefunctions and in the limit of very high density where the plasma is thought to have an extended (solid-like) structure, the molecular orbitals will most easily be classified into electronic bands. The techniques of solid-state theory will at this point most readily apply.

Conclusions

In this note, the transient molecules formed by collisions between plasma species have been investigated for the case of a hydrogen plasma containing H atoms and H^+ ions by using exact solutions of the Schrodinger equation for the H_2^+ molecule. Three major differences with the usual treatment of the electronic structure have been found

- (i) The two-centre treatment predicts the existence of absorption processes not accounted for in a single-centre description.
- (ii) The average oscillator strength of a transition that is allowed in the limit of large internuclear separation may be significantly altered in a two-centre description.
- (iii) A coupling between radiation and ionic kinetic energy occurs which has not previously been described.

Although only electric-dipole photon absorption processes have been investigated in this report, the same considerations could be applied to

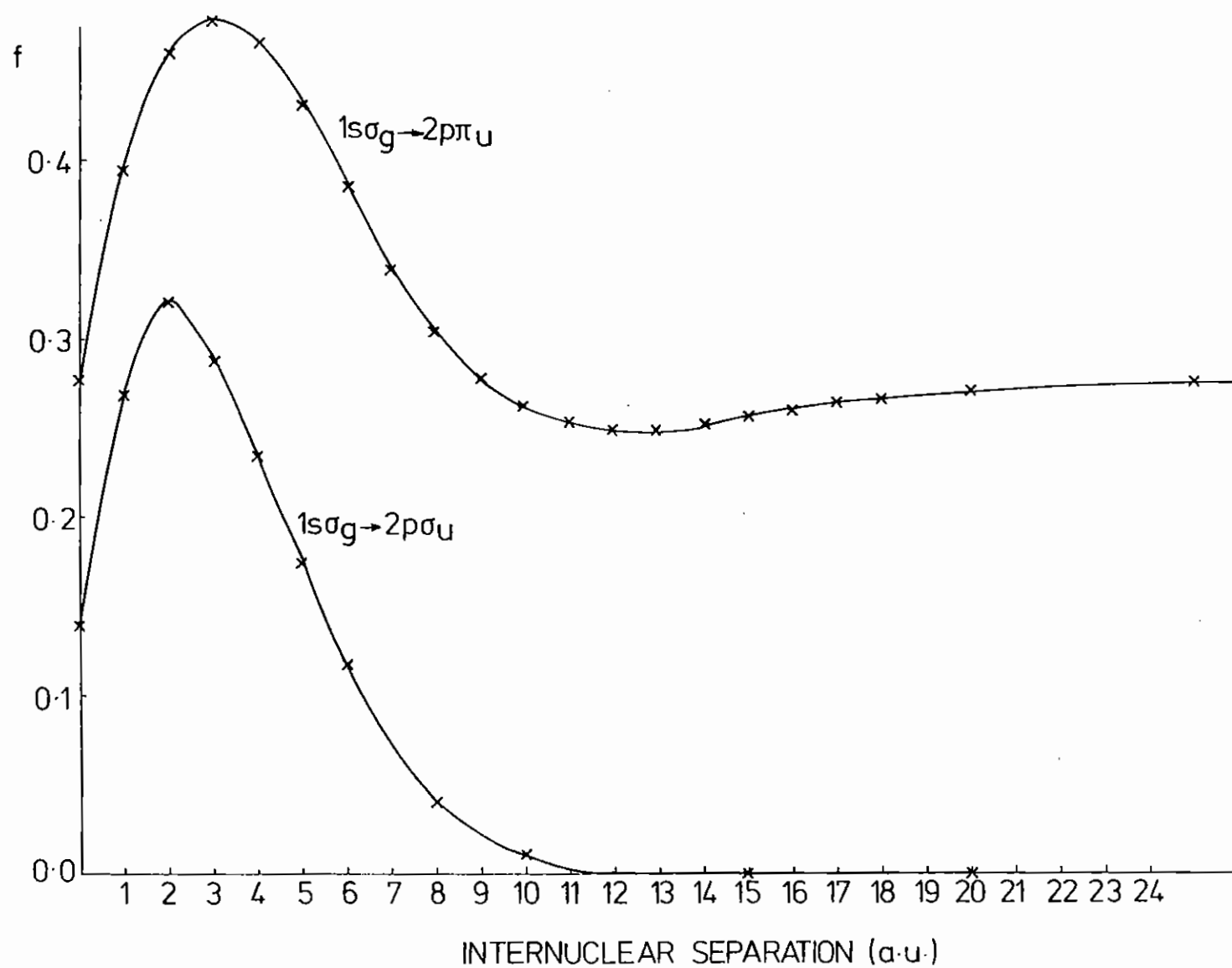


Fig. A5.28 The dependence of oscillator strength on internuclear separation $1s\sigma_g \rightarrow 2p\pi_u$ and $1s\sigma_g \rightarrow 2p\sigma_u$ in H_2^+ .

electron-impact excitation and ionisation processes. In discussing the ionic motion we have also assumed that the Born-Oppenheimer approximation holds and that the ions behave classically. In future work it is hoped to both provide more quantitative estimates of the effects predicted in this note and investigate the limitations of the assumptions made.

5.4.3 LTE Radiation transport in Hot, Dense high-Z and low-Z plasmas

T D Beynon and E H Smith (Birmingham)

As part of a program in ICF target physics a series of fast algorithms has been developed for use in the 1D hydro-code MEDUSA to compute LTE radiation transfer. Two methods have been optimized for use with low Z materials ($Z \leq 6$), where the plasma is usually fully ionized and the radiation mean free path is usually longer than the target size, and for use with high-Z materials where the opposite conditions obtain.

For the low-Z case, an energy-integrated semi-analytic solution to the integral transport is used with cross sections, $\Sigma(E)$, behaving as $\Sigma(E) \propto E^{-\gamma}$ where γ is an empirical constraint of approximate value 3. The application of integral transport techniques for planar geometry is relatively trivial, but there are no simple analytical representations for the concentric annular geometry found in ICF targets. Our approximation (A5.45) makes suitable modifications to the planar results to allow for the spherical case by considering the radiation flow to be made up of two components. The first 'direct' component is due to radiation which travelled inwards or outwards only. The second 'reflected' component is from radiation which initially travelled inwards but subsequently travelled outwards. The nett balance between the two is made in each MEDUSA mesh interval. A useful test of the complete algorithm is to try to calculate energy transfer in an effectively infinite sphere at a uniform temperature when no nett energy transfer should take place between each mesh interval. Over a wide range of opacities, the method is better than 1% in nett balance. More detailed checks were made using a discrete ordinate multigroup transport code ANISN (A5.46) in a spherical deuterium plasma. Figure A5.29 and A5.30 show the

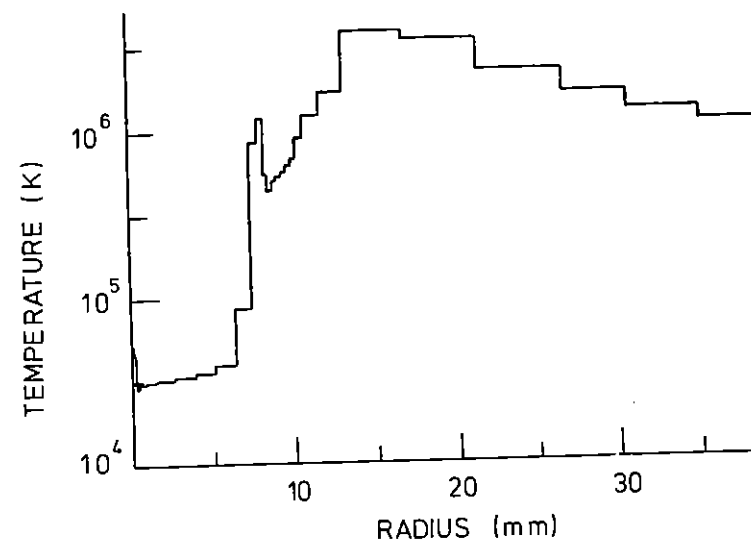


Fig. A5.29 Temperature distribution in the spherical deuterium plasma used for computational tests.

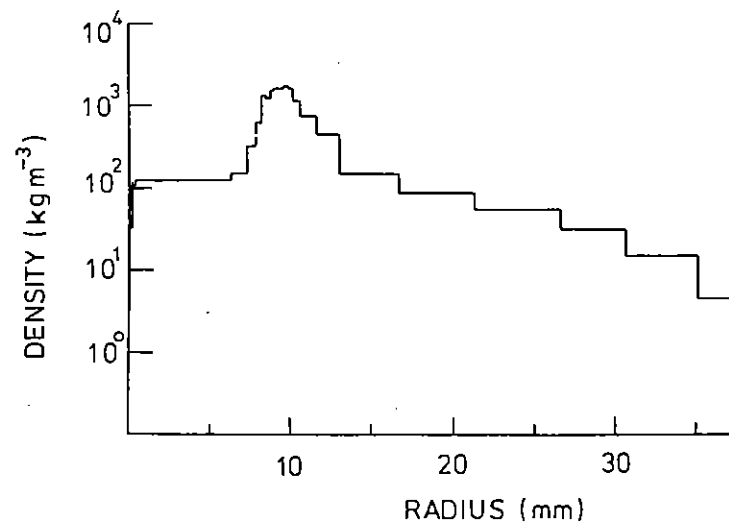


Fig. A5.30 Density distribution in the spherical deuterium plasma used for computational tests.

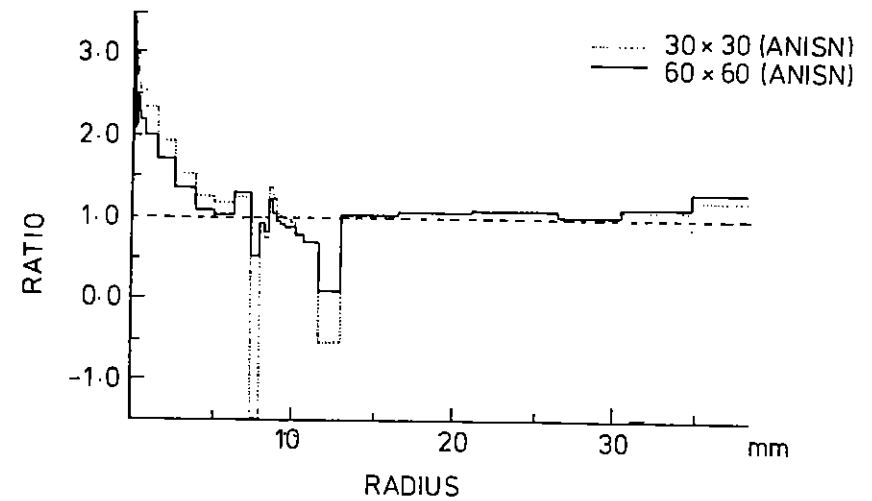


Fig. A5.31 Ratios of the computed energy leakages (ANISN/MEDUSA) for 30 energy group - 30 mesh interval and 60 - energy group - 60 mesh interval ANISN calculations in a spherical deuterium plasma.

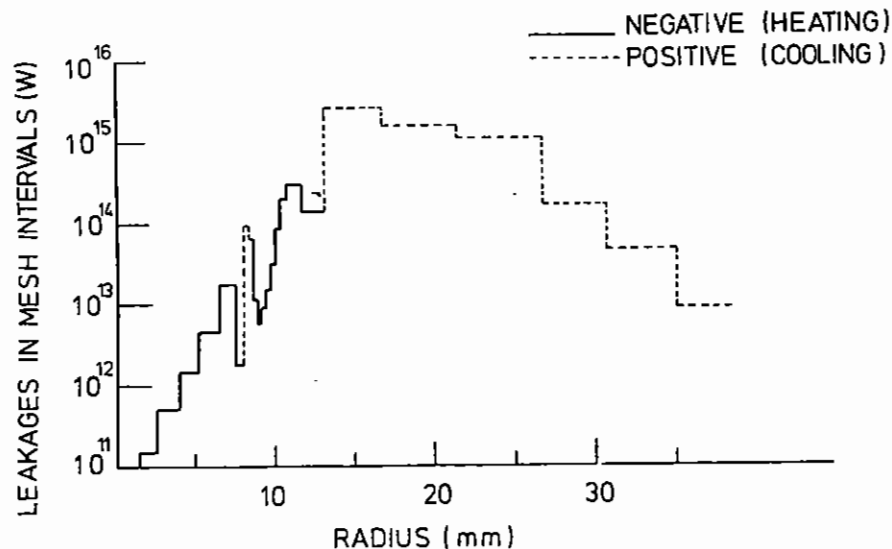


Fig. A5.32 Distribution of the net energy leakage rates from each mesh interval calculated by MEDUSA for the spherical deuterium plasma.

temperature and density distribution in the plasma as a function of radius. In figure A5.31 the energy leakages across each MEDUSA mesh interval are shown by each method and expressed as the ratio of the ANISN calculation to the present algorithm. Two ANISN runs are shown, one using 60 mesh intervals and 60 energy groups and the other using 30 intervals and 30 groups. An angular quadrature of S_{16} was used in each case. Fig. A5.32 shows the values of the energy leakages from each MEDUSA mesh interval calculated by the algorithm. All opacity data included bound-free and free-free cross sections. The execution time of the algorithm is of the order of 100ms per transport calculation, compared with about 20s for an ANISN calculation.

For the high-Z case, the radiation mean free path, λ_r , is usually considerably shorter than the size of any reasonable numerical mesh. Under these conditions the radiative heat flow behaves in a similar manner to electron heat conduction with very strong coupling between heat flow and the hydrodynamic behaviour. First order Taylor-expansion (A5.45) of the integral transport equations produces a Rosseland mean definition of K_r , an effective radiation conduction coefficient so that the energy current J , can be written as $J = -K_r \nabla T$ where T is the temperature distribution and $K_r = \frac{1}{3} \frac{\int_0^\infty \frac{dB(E)dE}{\lambda(E)}$ where

$$B(E,T) = \frac{E^3}{\pi^2 c^2 (\exp(E/kT)-1)}$$

is the blackbody limit flux. This definition of K_r is now incorporated within the MEDUSA treatment for the electron conduction via K_e .

Details of the treatment of consistent boundary conditions for using the high-Z and low-Z algorithms can be found in ref. A5.45.

A5.4.4 Determination of the opacities of low-Z elements for ICF targets.

T D Beynon, S Furnival and E H Smith (Birmingham)

A detailed calculation has been made of the opacity of low-Z elements ($Z \leq 6$) which are subject to a range of temperatures ($T = 1 \rightarrow 10^5 \text{ eV}$) and densities ($\rho = 10^{26} \rightarrow 10^{31} \text{ atoms/m}^3$) likely to be found in ICF targets.

An average atom model was used to predict the atomic properties necessary for cross-section calculations, energy levels and occupations of levels. The model consisted of either a density-modified Coulomb potential ($Z=1$) or Amaldi corrected Thomas-Fermi potential ($Z>1$) to describe the electron distribution. By summing the occupation of bound levels, the number of 'free' electrons could be determined. These were modelled using Fermi statistics to allow for the effects of degeneracy and were assigned a mean energy of $E \sim kT$.

Only two sources of radiation were considered, namely bound-free and free-free transitions. The bound-free cross-section was determined using a semi-empirical relationship due to Kramers (A5.47). For free-free transitions, it proves easiest to calculate the source and determine the cross-section from the principle of detailed balance the source being calculated using an equation due to Elwert (A5.48).

The cross-sections are converted to opacities by multiplying by the appropriate number of electrons in the initial and final states, the total opacity then being a sum of these two. Bound-bound transitions were ignored as the plasmas will be opaque and the highly discrete bound-bound transitions will be strongly re-absorbed.

Beynon and Smith (A5.49) have shown that in low-Z plasmas at temperatures greater than 100 eV, the opacity is effectively given by

$$\chi(E) \propto E^{-\gamma} \quad (1)$$

where E is the photon energy and $\gamma \approx 3$ and adjustable. To test the applicability of 1) to the range of ρ, T considered above, a dimensionless variable $D(E)$ was constructed.

$$D(E) = 8.4 \times 10^{-14} \epsilon(E) E^3 \frac{T^{3/2}}{n_e^2 Z^2} \quad (2)$$

where E and T are in eV and n_e in m^{-3} . $D(E)$ is consistent with the bremsstrahlung source term originally in MEDUSA. It was then averaged over all energy to get a single group averaged value for each ρ, T chosen.

The data, having been normalised, is plotted in Figure A5.33. In the region of high temperature ($T > 100 \text{ eV}$), equation 1) is a good approximation. At low temperature and low densities, the large peak is due to the dominance of bound-free transitions which go as n_e^{-1} whereas $D(E)$ varies as n_e^{-2} .

In the low temperature, high density region, there are two points of interest. One is the appearance of a bound-free edge, the other is the variation of $\langle D(E) \rangle$ with density for constant temperature, because of degeneracy.

A5.4.5 A semi-analytic radiative preheat model

C L S Lewis (Belfast) and S J Rose (RAL)

Introduction

A simple semi-analytic model is presented to describe the preheating effects in a shell target due to X-ray radiation from the ablation zone plasma which surrounds the target during laser irradiation. The model can aid in the design of targets to minimise preheating effects during ablative implosions. The model is 3-dimensional in the sense that the absorbed energy density from photo-absorption at any point in the shell is calculated by summing the contributions from all surface plasma elements when the spectral brightness of the surface plasma is known. The radiative preheat effect calculated from the model is compared to the effect calculated from the currently used version of Medusa containing radiation transport (A5.50). The effects of target scale size and aspect ratio are considered and the predictions of this 3-D model are compared to the predictions of an even simpler 1-D model.

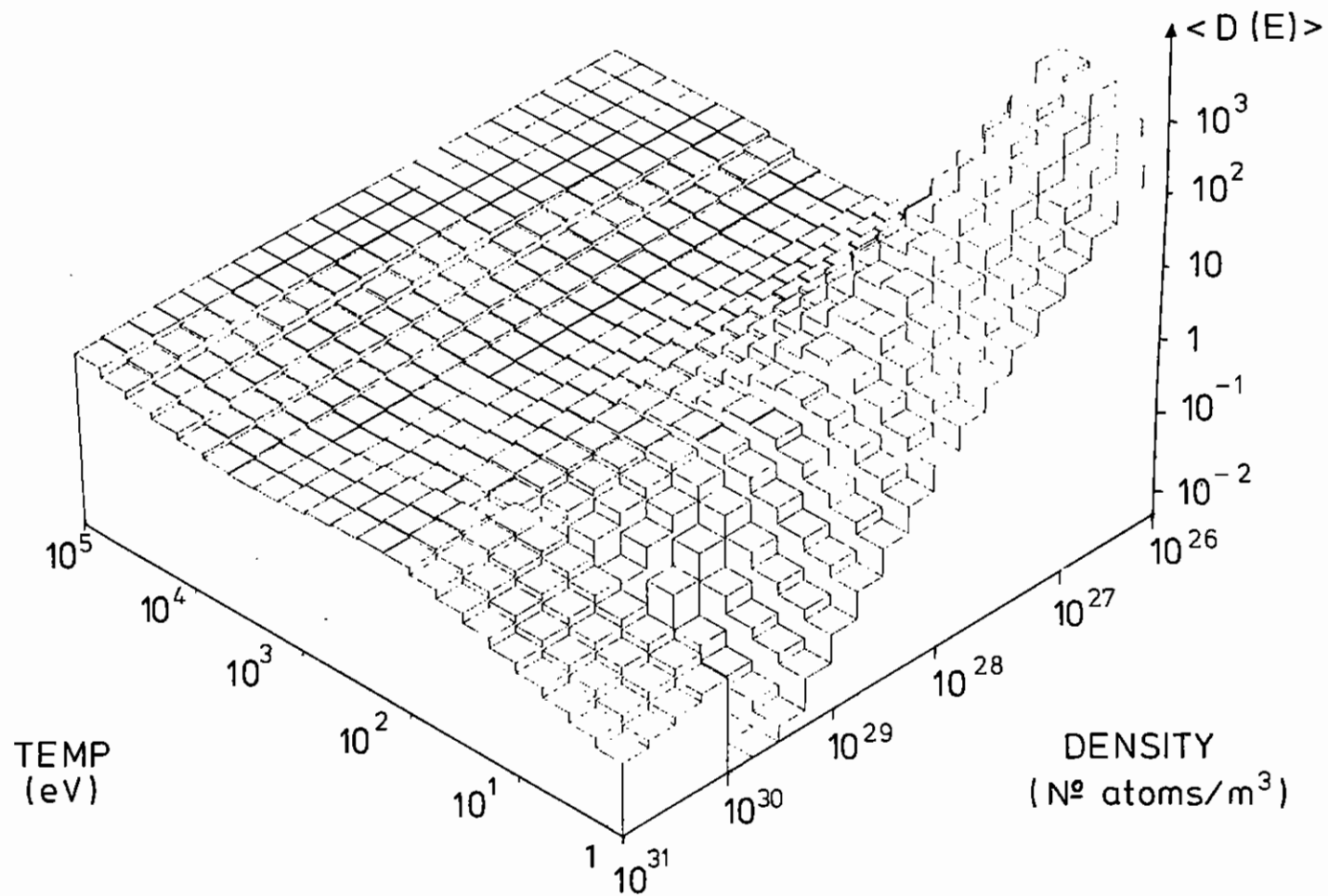


Fig. A5.33 A plot of the group averaged constant, $\langle D(E) \rangle$, for helium as a function of density and temperature.

Model

The geometry and symbols used in describing the model are illustrated in Figure A5.34. It is assumed that the relevant X-rays are emitted from a thin surface layer near the ablation front such that a $\ll x \ll \Delta R$, the surface emission brightness $I(E)$ in units of keV/keV/str/cm²/sec is measurable or known and that the outward flux equals the inward flux which causes preheat.

The intensity reaching the point P from a surface element of area ΔS at S due to X-rays of energy E to E + ΔE is given by

$$I_P = \frac{I(E)\Delta E\Delta S}{l^2} \exp[-\mu_s(l-m)] \exp[-\mu_g m \rho_g] \text{ keV/sec/cm}^2 \quad (1)$$

Where μ_s , μ_g and ρ_s , ρ_g are the mass absorption coefficients at photon energy E and densities of the shell material and fill gas respectively.

The energy absorbed at P between l at $l + \Delta l$ is

$$-\Delta I \rho = \mu_s \rho_s I_P \Delta l \quad (2)$$

and hence from eqn. 1 the absorbed energy density at P is given by

$$c_P = \frac{\Delta I}{\Delta l} = \frac{I(E)\Delta E\Delta S}{l^2} \mu_s \rho_s \exp[-\mu_s(l-m)] \exp[-\mu_g m \rho_g] \text{ keV/cm}^3/\text{sec} \quad (3)$$

Summing all ΔS elements about the axis of symmetry OP and all angles θ gives

$$\epsilon(x) = I(E)\Delta E 2\pi R^2 \mu_s \rho_s \int_0^\pi \frac{\exp[-\mu_s(l-m)] \exp[-\mu_g m \rho_g]}{l^2} d\theta \quad (4)$$

In the integration l and m are both functions of θ and x and simple geometry shows

$$l = (R^2 + (R-x)^2 - 2R(R-x)\cos\theta)^{1/2} \\ m = 2((R-\Delta R)^2 - R^2 + (R^2 - (R-x)^2 + l^2)^{1/2})^{1/2} \quad (5)$$

and $m = 0$ for $1 \leq (R^2 - (R-\Delta R)^2)^{1/2} + ((R-x)^2 - (R-\Delta R)^2)^{1/2}$
(This is the condition for rays tangential to the gas sphere).

Hence since $\sin\theta d\theta = dl/R(R-x)$ equation 4 can be written as

$$\epsilon(x) = I(E)\Delta E \frac{2\pi R \mu_s \rho_s}{(R-x)} \int_x^{2R-x} \frac{e^{-\mu_s \rho_s (l-m)} e^{-\mu_g \rho_g m}}{l} dl \quad (6)$$

with $m = f(l)$ as described in eqn. 5.

The total absorbed energy density is found by integrating over the entire source spectrum allowing for the energy dependence of the mass absorption coefficients. To simplify the integration it is assumed the effect of the fill gas is negligible i.e. $\mu_g \rho_g m \approx 0$ and that the shell material mass absorption coefficient can be represented as

$$\mu_s(E) = \mu^1/E^3 \text{ for } E > E_K \\ = \mu^1/JE^3 \text{ for } E < E_K \quad (7)$$

where μ^1 is the mass absorption coefficient (cm²/g) at $E = 1.0$ keV and J is the absorption coefficient jump ratio at the material K-absorption edge at $E = E_K$ keV.

The integral for the total absorbed energy density can now be broken into four parts representing the contributions above and below the K-edge which traverse the central void and those that traverse shell material only. i.e.

$$\epsilon(x) = \frac{2\pi R \rho_s \mu^1}{(R-x)} \int_0^{E_K} \frac{I(E)}{JE^3} \int_x^{l^1} \frac{\exp(-\mu^1 \rho_s l/JE^3)}{l} dl dE \\ + \frac{2\pi R \rho_s \mu^1}{(R-x)} \int_{E_K}^\infty \frac{I(E)}{JE^3} \int_{l^1}^{2R-x} \frac{\exp(-\mu^1 \rho_s (l-m)/JE^3)}{l} dl dE \\ + \frac{2\pi R \rho_s \mu^1}{(R-x)} \int_0^{E_K} \frac{I(E)}{E^3} \int_x^{l^1} \frac{\exp(-\mu^1 \rho_s l/E^3)}{l} dl dE$$

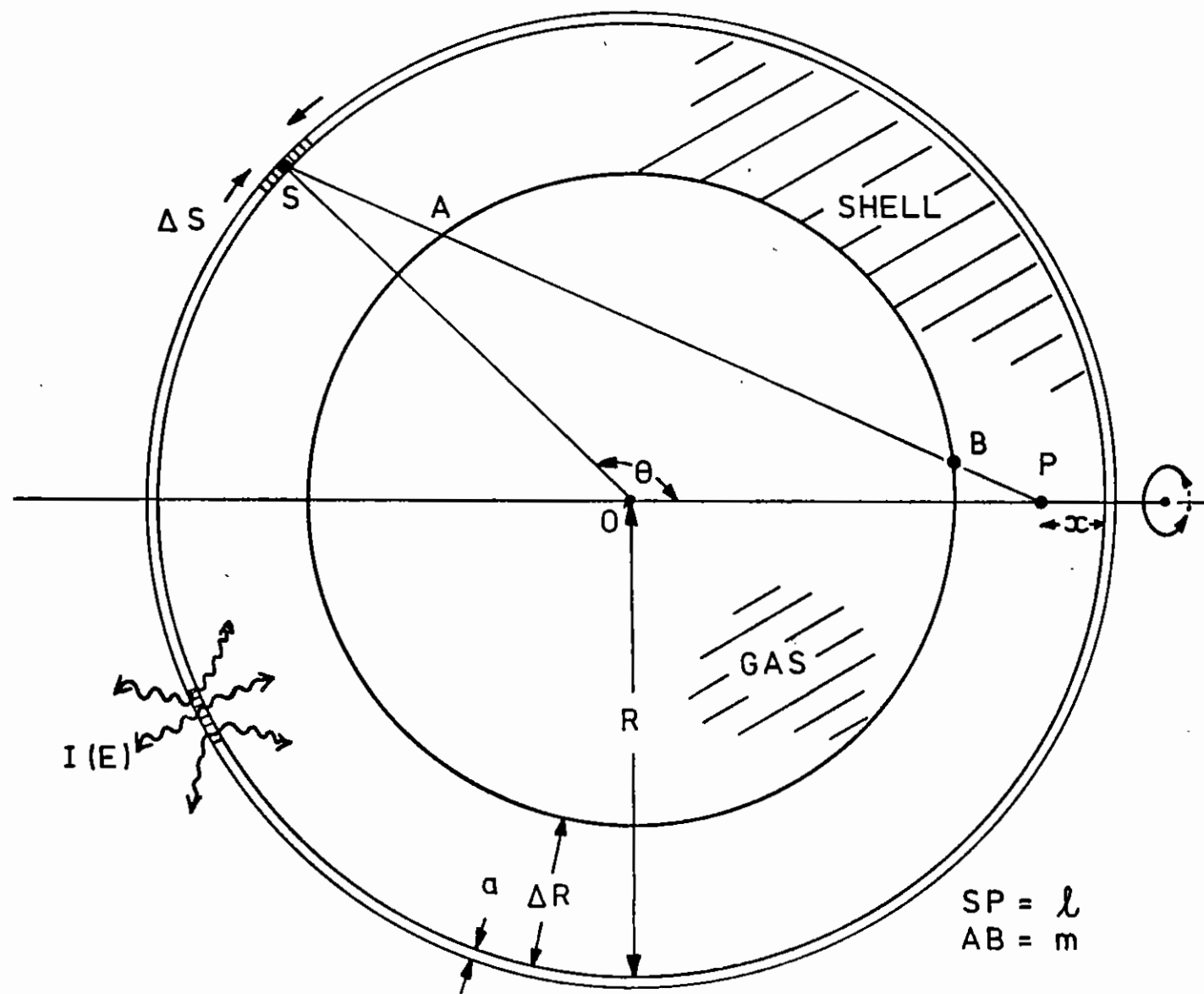


Fig. A5.34 Geometry and symbols used to evaluate the radiative preheat level at a point P.

$$\text{where } m = 2 \left[(R - \Delta R)^2 - R^2 + \frac{R^2 - (R - x)^2 + x^2}{2} \right]^{1/2} \quad (8)$$

$$l^1 = \left[R^2 - (R - \Delta R)^2 \right]^{1/2} + \left[(R - x)^2 - (R - \Delta R)^2 \right]^{1/2}$$

and $\Delta R > x > 0$ with $\epsilon(x)$ in units of $\text{keV}/\text{cm}^3/\text{sec}$

Application of Model

The integrals in eqn. 8 have been evaluated for the case of CH shell targets irradiated at $I_{\text{inc}} = 10^{14} \text{ W cm}^{-2}$ for 1 nsec at $\lambda = 0.53 \mu\text{m}$. The absolute emission brightness under these conditions has been measured experimentally and used to normalise the radiation transport treatment in Medusa (Ref. A5.50). The source function $I(E)$ used in the following is illustrated in Fig. A5.35 where a single temperature ($T_e = 200 \text{ eV}$) approximation to the Medusa calculated output continuum slope is taken. Cold material mass absorption coefficients are used in the present treatment (i.e. for CH, $\mu_1 = 2000 \text{ cm}^2 \text{ g}^{-1}$).

Figures A5.36 and A5.37 illustrate the preheat level dependence on the shell aspect ratio, $R/\Delta R$, with the surface laser intensity constant at $10^{14} \text{ W cm}^{-2}$. In Fig. A5.36 $\Delta R = 10 \mu\text{m}$ is held constant with R varied between $10 \mu\text{m}$ (solid sphere) and $\gg 100 \mu\text{m}$. Also shown is the curve calculated from a simpler model taking a 1-D flux of X-rays into the target. Here the absorbed energy at a depth x is expressed in $\text{keV}/\text{cm}^3/\text{sec}$ as

$$\epsilon(x) = 2\pi \rho_s \int_0^\infty \mu_g(E) I(E) \exp(-\mu_g(E) \rho_g x) dE$$

where all the symbols are as used in the 3-D model. The 1-D model has no R dependence and only ΔR dependence in setting the maximum useful value of x . For shells with aspect ratio $R/\Delta R \geq 4$ the preheat distribution within the shell is almost constant but overestimated in this case by the 1-D model by a factor of ≈ 2 . In general this discrepancy factor depends on the source plasma temperature (T) and the material mass absorption coefficient (μ_g).

Figure A5.37 illustrates the preheat effects when the aspect ratio is varied by changing the shell thickness. As the shell is made thinner, with a constant radius R the midshell preheat level increases with the inside surface level increasing at an even faster rate due to radiation from the other side of the target. The 1-D model prediction at a given depth in the

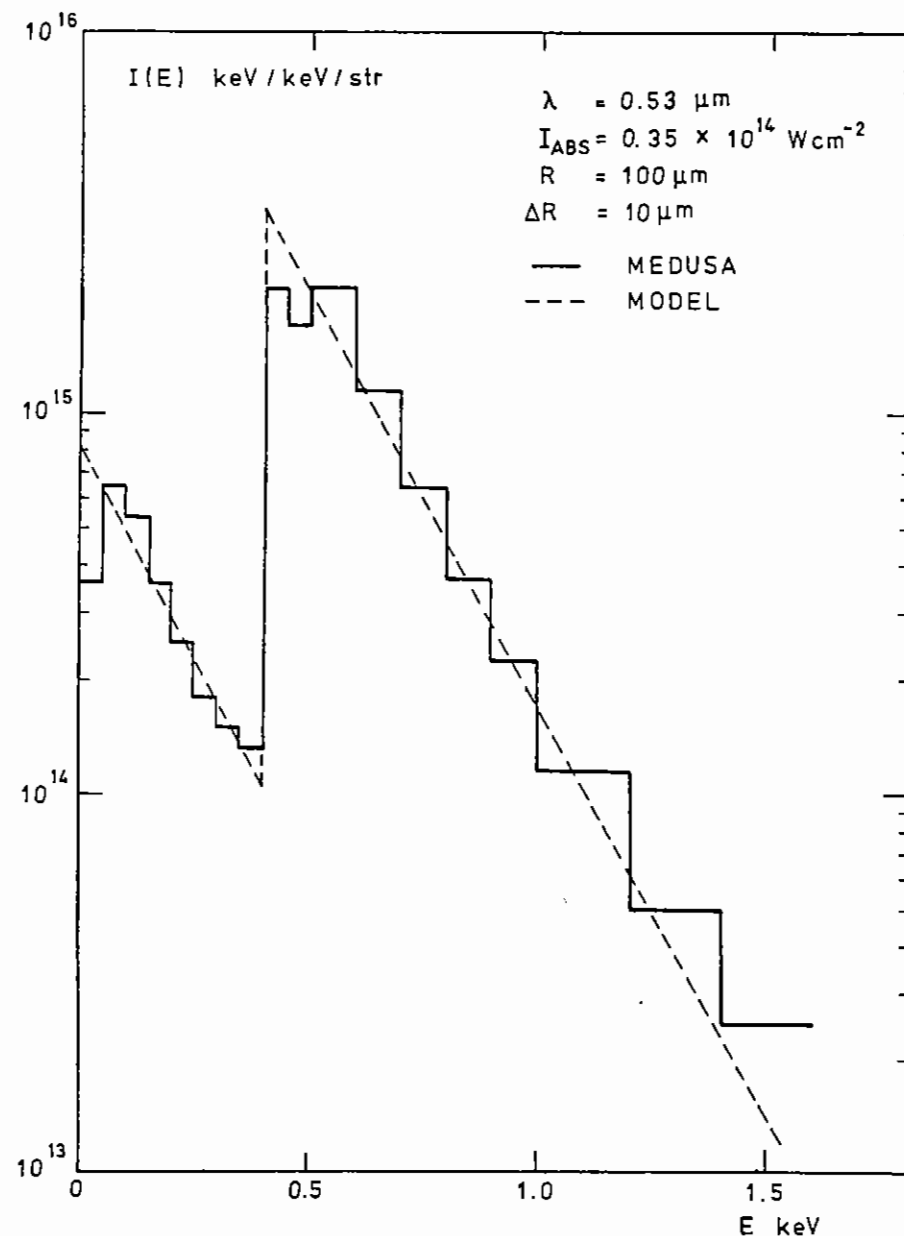


Fig. A5.35 The source spectrum used in present calculations of preheat.

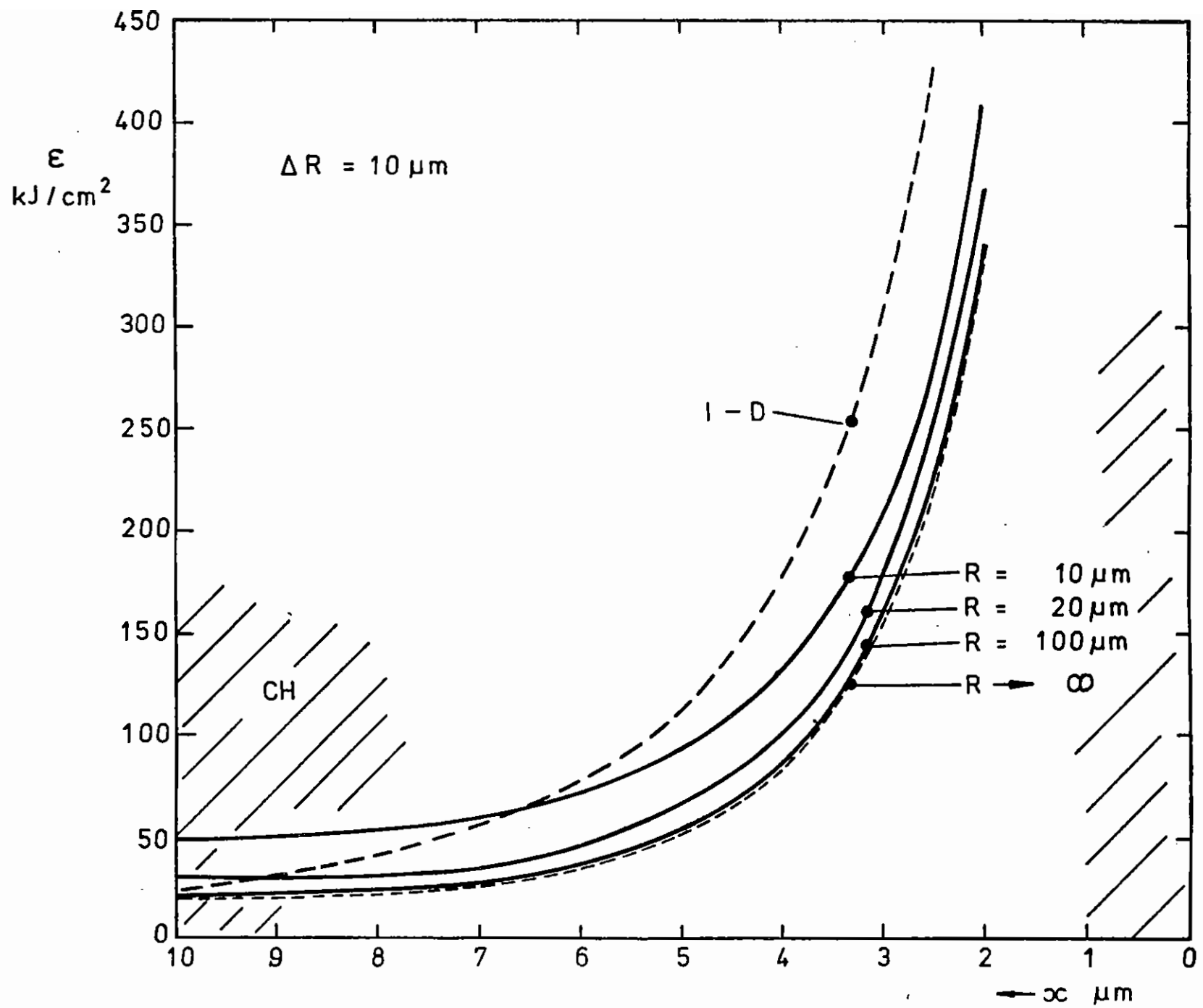


Fig. A5.36 Preheat dependence on aspect ratio when R is varied.

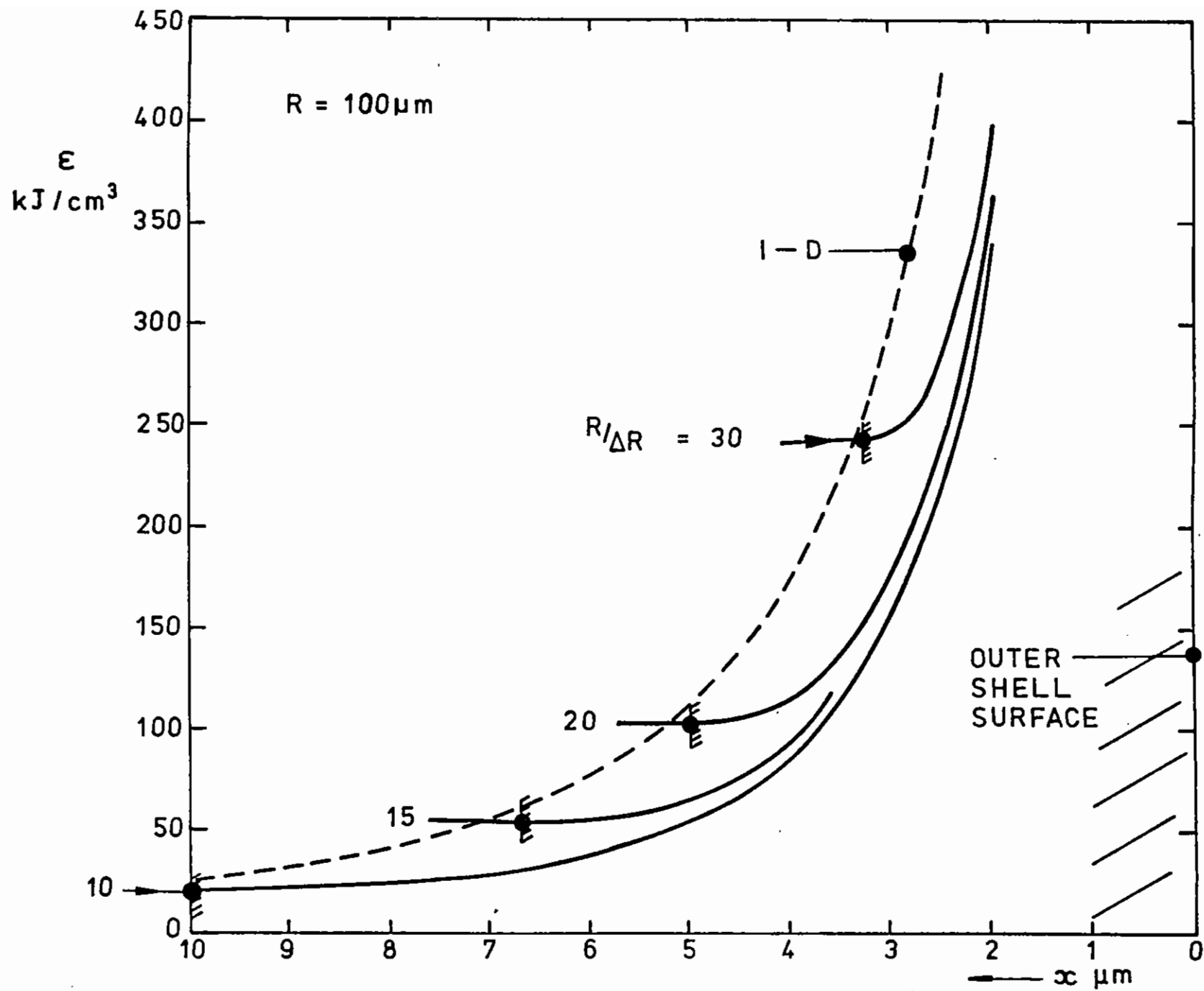


Fig. A5.37 Preheat dependence on aspect ratio when ΔR is varied.

shell makes no distinction between shells of different wall thickness.

Figure A5.38 compares the models with a Medusa simulation for the same target and irradiance parameters. Values of the Medusa energy density are obtained using the calculated temperatures together with the equation of state and therefore include shock heating and hydrodynamic effects although an attempt has been made to minimize these by running the laser pulse for a fraction of its Insec duration and scaling to Insec. The dominant radiative effects agree well with the 1-D model prediction. The model calculates preheat in the form of absorbed energy density per unit time and Figure A5.38 indicates the material temperature equivalent using an equation of state for carbon at $\rho = 1 \text{ g cm}^{-2}$ and a heating duration of 1 nsec.

The model allows for a quick assessment of the relative importance of accurately knowing the shell material mass absorption coefficient. This is relevant to Medusa code development where the radiation transport treatment will include atomic physics and ionisation fractions. Figure A5.39 illustrates one aspect of this where the absorption edge position has been varied from $E_K = 0.28 \text{ keV}$ for cold material to $E_K = 0.39 \text{ keV}$ the ionisation potential of He-like carbon. For the same source function $I(E)$ as before the effect at the inside surface of the shell is insensitive to the exact detail of this energy window.

Some general features of the 3-D model analytic expressions can be deduced. The preheat level at the centre of a solid sphere is always exactly double that predicted by the 1-D model. This is seen from equation 8 with $R = \Delta R$ and $x = \Delta R - \sigma$ in the limit $\sigma \rightarrow 0$. The preheat level at the inside surface of any shell is calculated from the two surface components coming from either side of the plane tangential to the gas sphere. These two contributions are identically equal and this can be verified from equation 8 with $x = \Delta R$, $\ell^1 = (R^2 - (R - \Delta R)^2)^{1/2}$, $m = (\ell^2 - (\ell^1)^2)/\ell$ and using the substitution $L = (\ell^1)^2/\ell$ in the integrals which include the void region. The half-plane contributions equal the 1-D prediction in the limit where $\Delta R = R$ as above.

A5.5 Laser Physics

A5.5.1 Transfer of Pump Intensity Variations to a Stokes Seed

S Jorna (St Andrews)

The steady-state coupled equations for the pump and Stokes fields are solved in the paraxial approximation to study the effect of phase and amplitude variations in the pump on the Stokes beam. Allowance is made for pump depletion and diffraction.

Coupled equations

In the slowly-varying amplitude approximation, the equation describing the evolution of the vibrational coordinate, Q , of the Stokes field, E_{si} , and the i -th mode of the pump field, E_{pi} , are (A5.51).

$$\frac{\partial Q}{\partial t} + \frac{1}{2} \Gamma Q = iK_1 \int_1 E_{pi}^* E_{si} e^{-i\Delta k_1 z}, \quad (1)$$

$$\frac{n_p}{c} \frac{\partial E_{pi}}{\partial t} + \frac{\partial E_{pi}}{\partial z} + \frac{1}{2ik_{pi}} \frac{1}{r} \frac{\partial}{\partial r} \left(r \frac{\partial E_{pi}}{\partial r} \right) = -i \frac{\omega_p n_s}{\omega_s n_p} K_2 Q E_{si} e^{-i\Delta k_1 z}, \quad (2)$$

$$\frac{n_s}{c} \frac{\partial E_{si}}{\partial t} + \frac{\partial E_{si}}{\partial z} + \frac{1}{2ik_{si}} \frac{1}{r} \frac{\partial}{\partial r} \left(r \frac{\partial E_{si}}{\partial r} \right) = iK_2 Q^* E_{pi} e^{i\Delta k_1 z}, \quad (3)$$

where Γ is equal to the FWHM Raman line width, $\Delta k_1 = k_{pi} - k_{si} - k_r$ is the wavenumber mismatch of the i -th mode, ω_p (ω_s) the pump (Stokes) frequency, and n_p (n_s) the refractive index at the pump (Stokes) frequency.

In the steady-state regime these equations may be combined to yield for a single mode.

$$\frac{\partial E_p}{\partial z} + \frac{1}{2ik_p} \frac{1}{r} \frac{\partial}{\partial r} \left(r \frac{\partial E_p}{\partial r} \right) = -\frac{1}{2} g_s \frac{\omega_p}{\omega_s} \left| E_s \right|^2 E_p \quad (4)$$

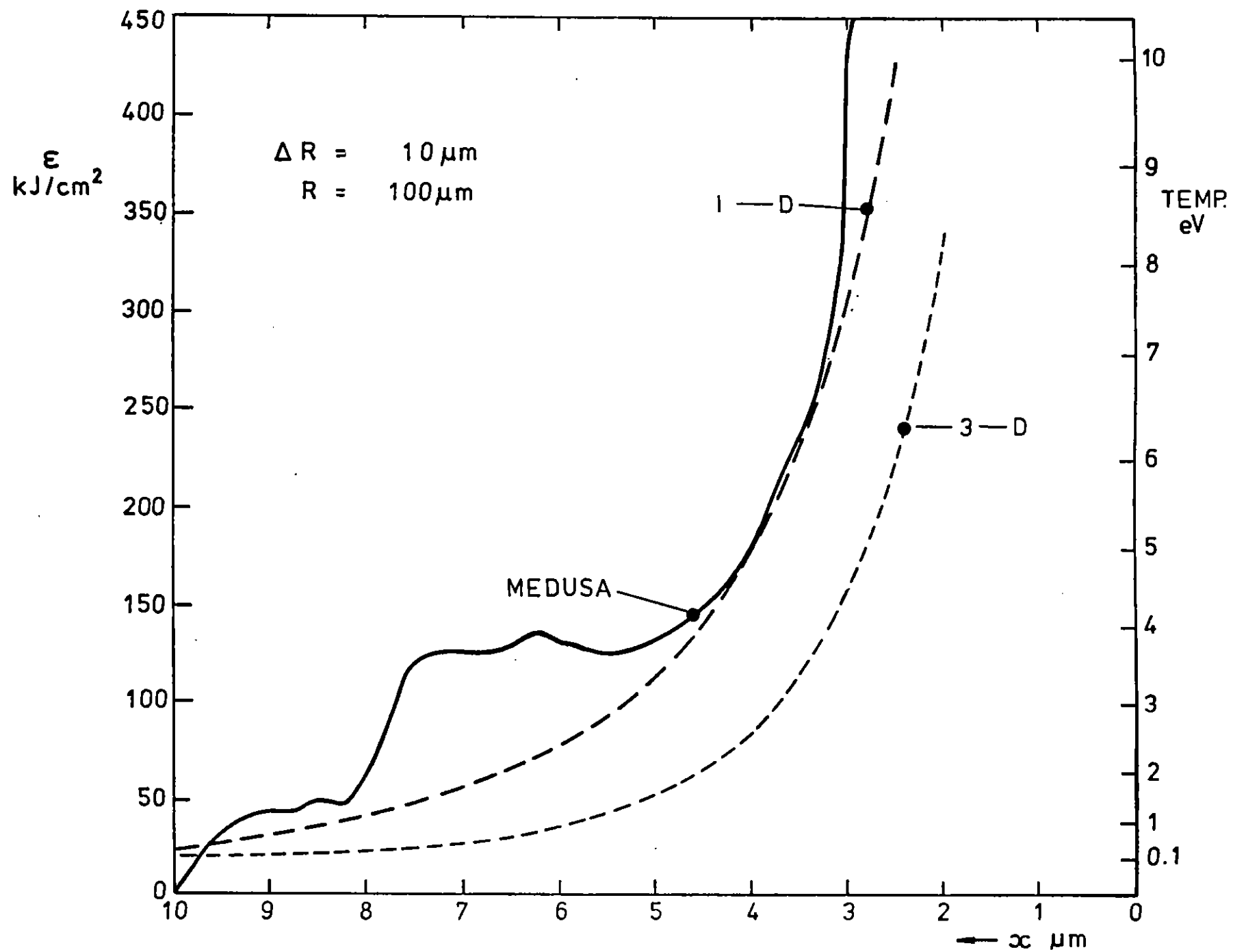


Fig. A5.38 Comparison between model and MEDUSA predictions.

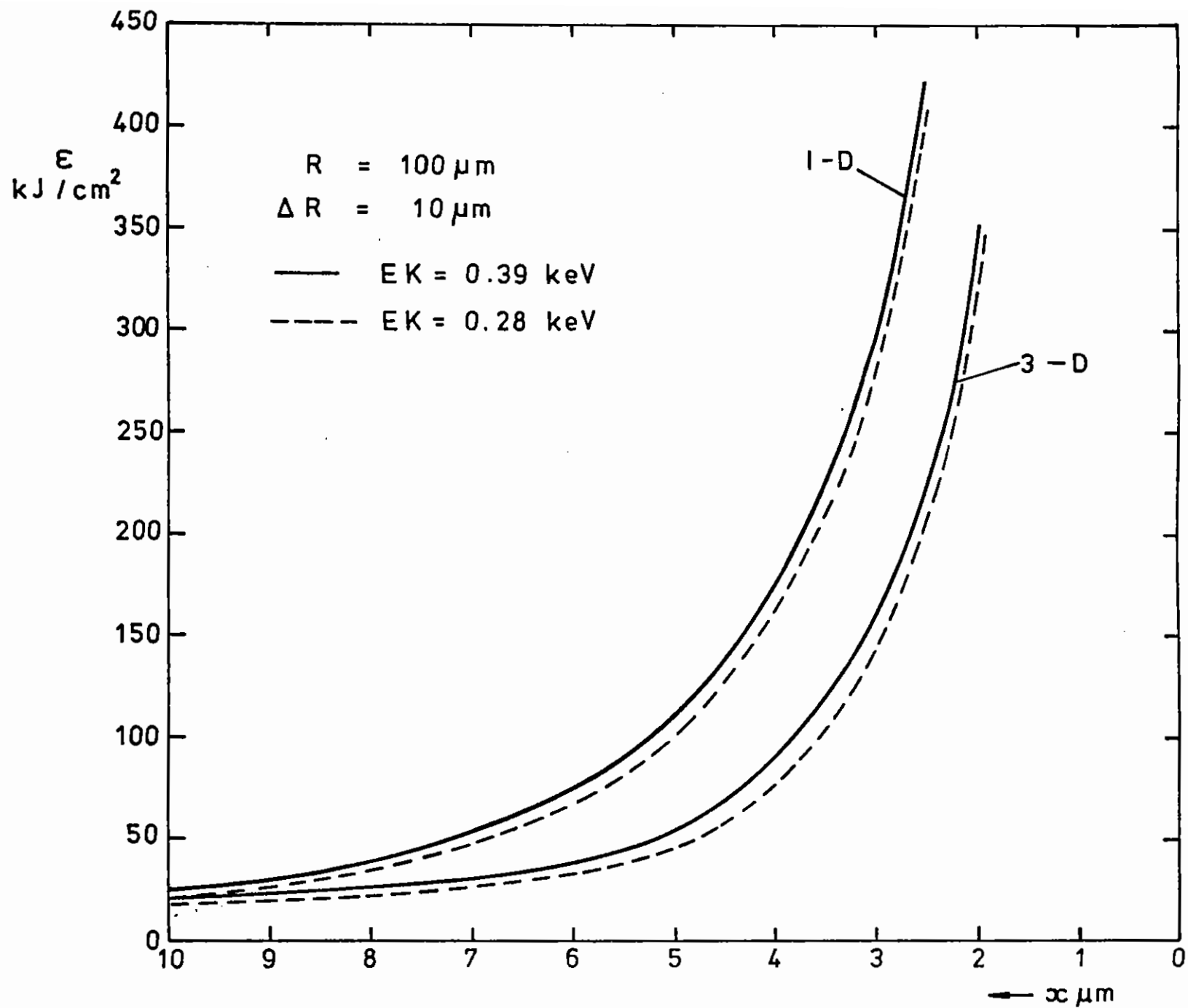


Fig. A5.39 Preheat sensitivity to nature of material mass absorption coefficient energy dependence.

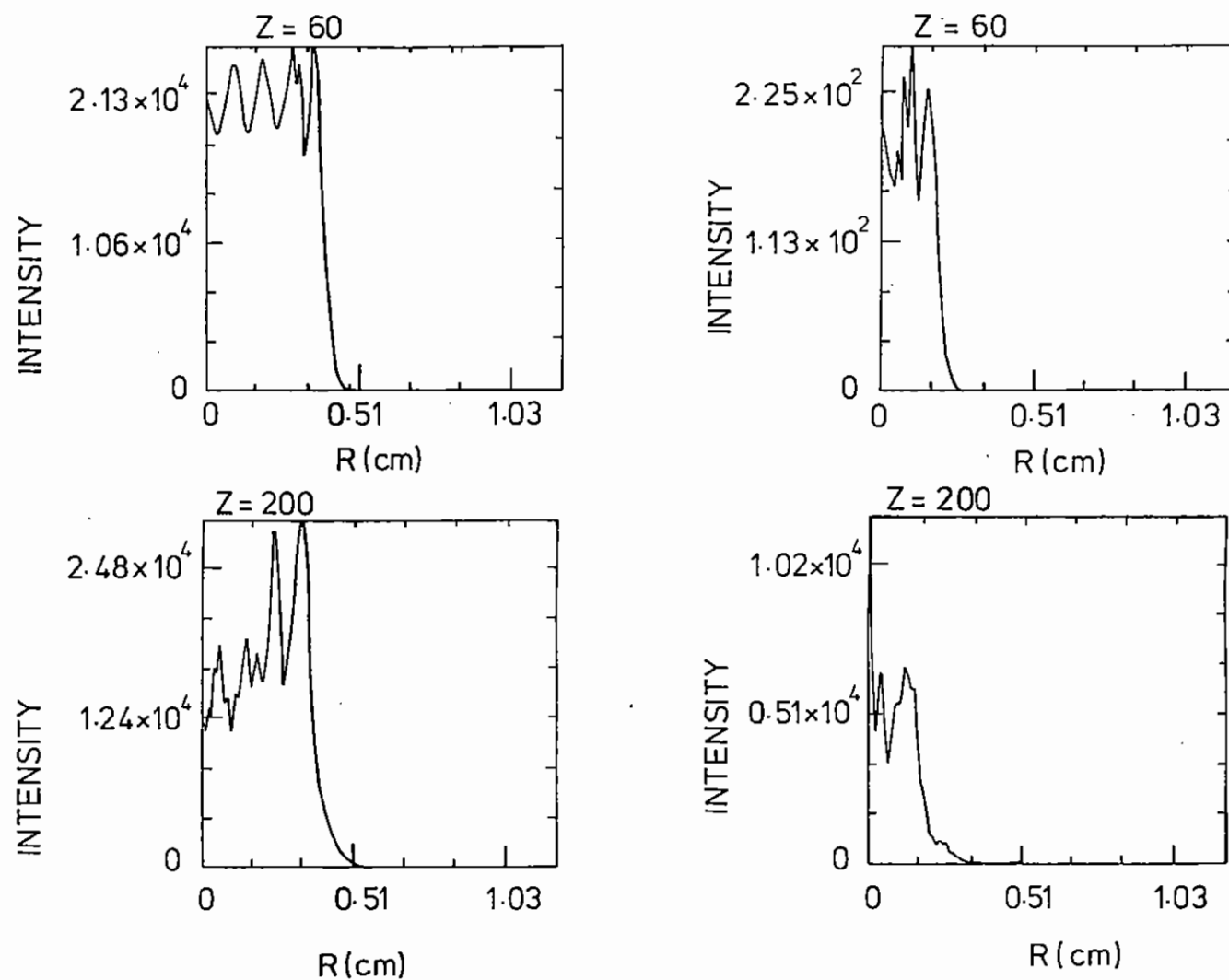


Fig. A5.40 Radial variation of pump (EL) and Stokes (ES) intensities at 60 cm and 200 cm into the Raman cell. Initial pump modulation was taken as $0.1 \cos(8\pi r/r_{\max})$ with $r_{\max} = 0.8$ cm.

$$\frac{\partial E_s}{\partial z} + \frac{1}{2ik_s} \frac{1}{r} \frac{\partial}{\partial r} \left(r \frac{\partial E_s}{\partial r} \right) = \frac{1}{2} g_s |E_p|^2 E_s \quad (5)$$

where g_s is the single mode Stokes power gain. These equations have been solved numerically by a scheme that allows implicit or explicit advancing in z and in which the r and z meshes are variable to optimize computational efficiency.

As an application we have chosen Raman conversion in hydrogen gas at 10 atm. We are well into the pressure broadening regime where the Stokes gain coefficient is independent of pressure. The pump diameter is 0.8 cm and the Stokes beam diameter is 0.4 cm. Both beams have a square intensity profile with a Lorentzian fall-off. The pump power flux is 10 MW/cm² propagating collinearly with a Stokes beam of initial intensity 0.02 MW/cm². The pump has an initial intensity modulation of 10% and a constant phase. Also, $k_p = 1.257 \times 10^5 \text{ cm}^{-1}$, $k_s = 9.956 \times 10^4 \text{ cm}^{-1}$, and $g_s = 2.4 \times 10^{-9} \text{ cm/W}$. The figures show the radial variation of the pump (EL) and Stokes beam (ES) intensities at 60 cm and the end of a cell of length 200 cm. The pump intensity was modulated by 10% at the beginning of the cell. The transfer of the spatial variations to the Stokes wave is clear from Figures A5.40b and A5.40d.

Conclusions

The results show that an initial intensity ripple on the pump is transferred to the Stokes beam which is further affected by the edge diffraction pattern propagating inwards. because the Stokes field is driven by the pump intensity, phase variations tend to be smoothed out.

The far-field calculation shows that when the Stokes energy is normalized to its value in the absence of gain, energy loss is less than 10% for this case.

Edge effects could be ameliorated by apodization with a decrease in conversion efficiency.

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A6 GLASS LASER FACILITY OPERATIONS AND DEVELOPMENT

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A6 GLASS LASER FACILITY OPERATIONS AND DEVELOPMENT

A6.1 Introduction

J E Boon and M H Key (RAL)

Facility operations have been characterised by reliable and intensive operation of the Nd glass laser (VULCAN) and its associated six beam and two beam target irradiation laboratories (TA1 and TA2). The laser has fired 1900 full power shots for scheduled experiments and the two target areas have been operated in parallel with alternate shots being directed to each area.

Facility development has centred on the new Phase I programme approved by the Science Board in its 1983 Forward Look and given Government (DES) financial authorisation in August 1983. This programme is associated with a major proposal document submitted by scientific users of the CLF detailing "Proposals for the long term scientific programme and related new developments at the CLF".

The Phase I programme includes development of both the Nd glass laser and of the high power KrF laser (The latter to 100 J, 1 nsec operation, see Section C).

The development of the glass laser will proceed in stages. First it will be provided with two improved target irradiation facilities. For more uniform irradiation of spherical targets 12 beam dodecahedral irradiation will be provided in Target Area West (TAW), together with a powerful short pulse diagnostic beam. Six beam high energy irradiation of plane targets, six beam line focus irradiation for X-ray laser research and a powerful single beam for high intensity interaction studies will be provided in Target Area East (TAE).

The power of the single beam will be increased in two stages, first to 400 J, 1 ns by provision of a 150 mm amplifier, then by a further power doubling to give up to 800 J, 1 ns.

The power of the multibeam laser will be doubled to give 2.4 kJ, 1 ns by increasing the number of 108 mm output disc amplifiers to 12 and modifying the feeder system.

The timescale for providing the new facilities is constrained both by available manpower and by the rate of funding. The latter is particularly critical. The CLF is fully operational and its allocation of funds includes typically only a 15-20% margin for new capital expenditure. Thus relatively small variations in the allocation can cause major perturbations to new capital projects. In particular a 6% reduction in allocation relative to the Forward Look guideline removed 35% of the new capital money in the reporting year. On current guidelines we should be able to report the completion of Phase I in our Annual Report for or 1988. Progress this year has centred on the construction of TAW (A6.5), which will start operation in June 1984. In preparation for this the six beam target irradiation lab (TA1) was shut down at the end of April after five years of valuable experimental research. Construction of TAE has been delayed to next year by the budget cut. The 150 mm amplifier is also well advanced with a prototype constructed during the year.

Aside from these major activities there has been steady progress on minor improvement and diagnostic developments which are also described in the following sections.

A6.2 Operations

A6.2.1 Glass Laser Operations

N Allen, C Danson, C Edwards, B Espey, A Raven, P Rumsby, R Wyatt, D Smith, S Sunderland and R Wellstood (RAL)

The laser system has performed very reliably during the whole of the 12 month period since the completion of the 1kJ enhancement programme in February 1983. The many major changes made to the layout of the whole system during this upgrade (detailed in last years report) have led to

significant improvements in system versatility and ease of operation and maintenance. Of the 52 weeks in the year 40 were dedicated to target area shots and the remainder used for maintenance and development of the system.

Once again laser availability has been significantly improved and during experimental periods the system now runs for 45 to 50 hours per week. Performance has been limited only by target area requirements with energies up to 1kJ in 1ns pulses and powers up to 4TW in 70ps pulses being delivered in the 6 beam compression facility. This area (TA1) received 990 shots during the year while the other target area (TA2), now with 2 beams available, received over 900 shots. A further 650 shots have been fired for test and development purposes. The shot failure rate due to all causes was less than 3% over the year.

Considerable on and off line development work has been performed during the year. Improvements have been made to beam diagnostics, oscillator reliability and energy monitoring. A new 150mm diameter disc amplifier has been constructed and the efficiency of the existing 108mm amplifiers examined in detail. A test facility for small flash lamps has been built and a damage test facility set up. These items are reported in detail in section A6.3.

A6.2.2 Target Area Operations

D Bassett, A J Cole, R Eason, C J Hooker, S Knight, D Pepler, W T Toner (RAL)

TA1 and TA2 operated in tandem throughout the year apart from a ten week shut down of TA2 to complete the link to the new areas (Section A6.5). Target Area East is currently used as an extension to TA2 and houses the Z-pinch (Section A1.6). Two beam operation of TA2 has been very successful and has led to increased user demands for time. The only significant problem has been an increase in the frequency of instances where the high repetition rate of shots in the double passed disc amplifiers (because TA1 and TA2 are using the same beams) leads to noticeable deterioration in beam quality.

Improvements in the laser chain have made higher output energies routinely available in both areas. This led to some problems of damage in the windows of the harmonic conversion crystals which were attributed to acoustic shocks in the index matching fluid due either to absorption by particles in the fluid or to non-linear processes. The Cargille 5040 fluid used has now been replaced with Koolase which work at Rochester and Livermore suggests is more tolerant to high intensities.

Most of the target area staff are now trained to work in either area in order to give increased flexibility in scheduling, particularly at holiday times, and also to enable them to carry out extended research and development tasks without disruption to routine operations.

TA1 fired its last shot on April 19th 1984. Short pulses in the three east side beams were used to drive an exploding pusher implosion to produce alpha particles to backlight a target driven in long pulse by the three west side beams in order to study the growth of the Rayleigh Taylor instability.

A6.3 Glass Laser Development

A6.3.1 Oscillator development

C Danson, C Edwards, A Raven, D Smith and R Wyatt (RAL)

Single Longitudinal Mode Resonator Stabilization

During this year a project was completed to build a feedback system to stabilize the optical resonator length of the single longitudinal mode oscillator. Oscillators of this type suffer from thermal and mechanical drifts, disturbing the optical path length of the cavity, leading to multi-mode generation. Beating then occurs between the modes producing a temporally modulated pulse output. The feedback system is designed to sample pulses produced during the countdown for a main laser shot. When 'beating' is detected the cavity length is adjusted by changing the potential applied to a piezo-electric crystal driven mirror.

A block diagram of the feedback system is shown in Figure A6.1. The signal from a pin diode, monitoring the 25ns wide oscillator output, is passed through a high pass filter, which discriminates between the required natural envelope of the single mode, and the higher frequency beating which may occur. The design requirements of the filter were that it should have a transmission of - 3dB at 40MHz and 60dB at 20MHz. The output of the filter is passed on to an amplifier and rectifier circuit, producing a d.c. level which is fed into a comparator. An output signal from the comparator triggers a series of up/down counters, which count on a pre-selected number of pulses from a 50 Hz pulse train. Selection of the number of pulses (eventually output voltage steps) is made using a hexadecimal coded switch, giving a 1 to 16 pulse range. A digital to analogue converter follows the counters and its output is passed on via amplifiers to a piezo-electric translator. The digital to analogue converter has a range of 257 steps. At this number and zero the counters are reversed, so providing continuous operation up and down through the voltage range at the chosen resolution. The piezo crystal movement range is about 8µm for a change of 1KV. The resolution can be varied by changing the overall voltage used since the 257 steps used are constant.

Several tests were carried out using different step sizes to judge effects on stability. Figure A6.2 shows the results of one test with a step size of 60V, i.e.: a resolution of about 0.5µm. A step on the trace occurs each time mode beating is detected. In Figure A6.2(a) the piezo-electric mirror was disconnected from the feedback system, whereas in Figure A6.2(b) the feedback loop was completed. In the case where the feedback system was completed, mode-beating occurs less frequently by a factor of 2, compared to a free running oscillator. In all the tests carried out, no advantage could be found for using one step size rather than another. This work was carried out on an off line single longitudinal mode YAG oscillator. The system is now being installed on the main laser YLF oscillator system.

A6.3.2 108mm Disc Amplifier Testing

N Allen, J Boon and C Danson (RAL)

The 108mm disc amplifiers each contain twenty lamps in a cylindrical geometry around the discs. The support members for the discs encroach into the lamp housing, partly obscuring four of the lamps from the discs. A programme was carried out to determine the effect of running the amplifiers with these lamps removed and whether the reduction in gain that would arise could be recovered completely by design modifications to the chassis and reflector housings. These modifications are demonstrated in Figure A6.03 which also shows an unmodified chassis. Running amplifiers with sixteen lamps only would make significant cost savings to any future upgrade scheme.

Two experimental periods were used for tests. In both cases TAl experiments required the use of four beams only of the six beam system, enabling the disc testing programme to use the remaining beam lines.

An absorbing glass calorimeter was placed at the output of each beam. One was used as a standard to measure the input to the system, while the other was used to calculate the amplifier gain. A schematic of the experimental layout is shown in Figure A6.4. As a cross check to the measurements the computer linked calorimeters used for normal shot energy measurements, were also used. Gain measurements were calculated using two shots, one with the test disc amplifier passive the other active. The standard disc measurements enabled allowances to be made for input variations. Most gain measurements were calculated using four shots, but if any inconsistencies arose further shots would be fired. Gain coefficients were calculated from the input and output energy data (assuming no saturation effects).

Initial tests looked very promising. Pairs of lamps were disconnected from an ordinary twenty lamp amplifier and measurements were taken to find the drop in gain. From this the contribution of each lamp to the overall gain could then be calculated. Most of the lamps in the amplifier were found to contribute 5-6% each to the gain coefficient (as expected) whereas the lamps

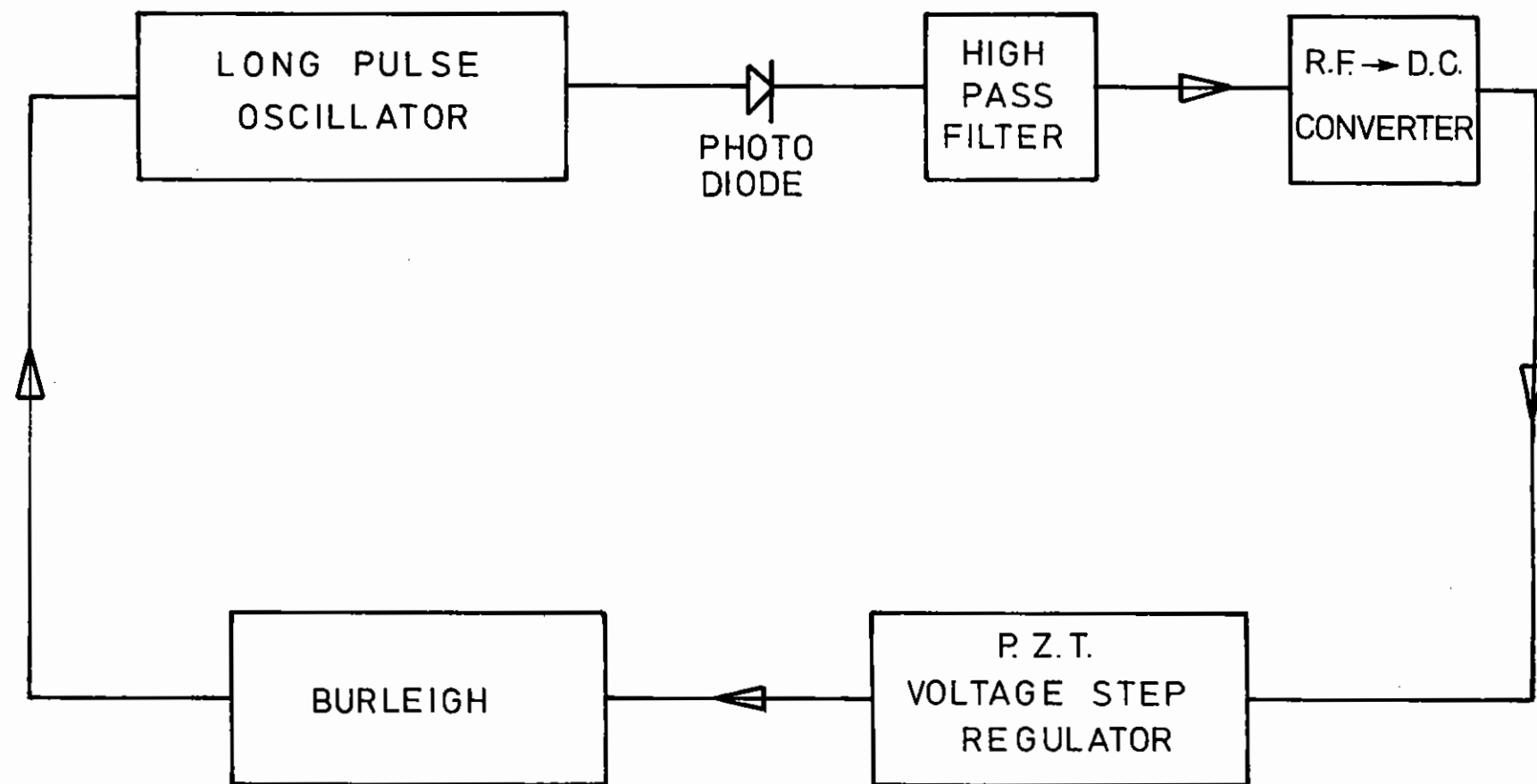


Fig. A6.1 Block diagram of the control for the single longitudinal mode oscillator rear mirror drive.

TESTING TIME $\approx$ 100 mins.



(a) FEEDBACK CONTROL DISCONNECTED.

(b) FEEDBACK CONTROL CONNECTED

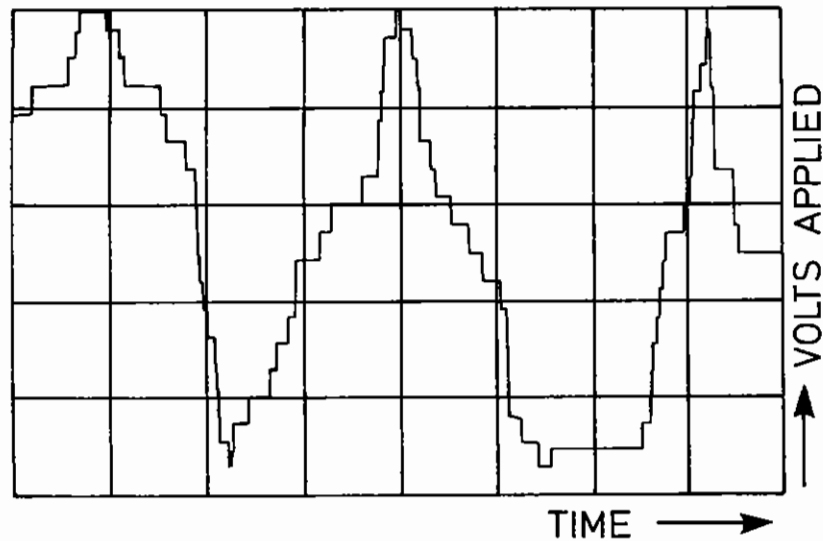
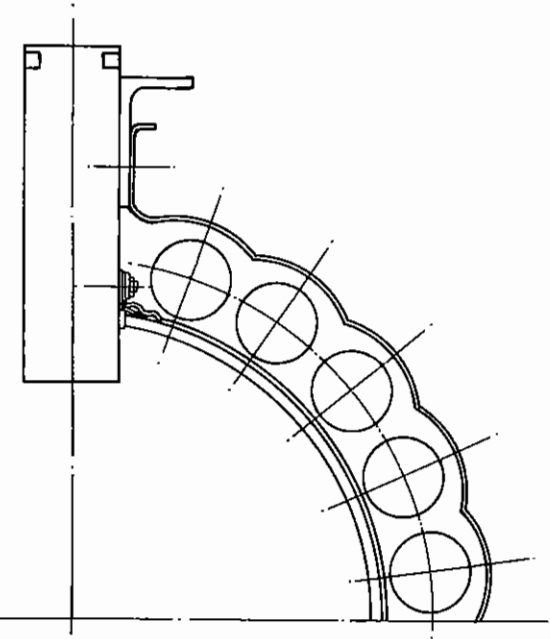


Fig. A6.2 Results of tests on stability of piezo electric controlled mirror.

(a)

20 LAMP
UNMODIFIED
CHASSIS



(b)

16 LAMP
MODIFIED
CHASSIS AND
REFLECTOR

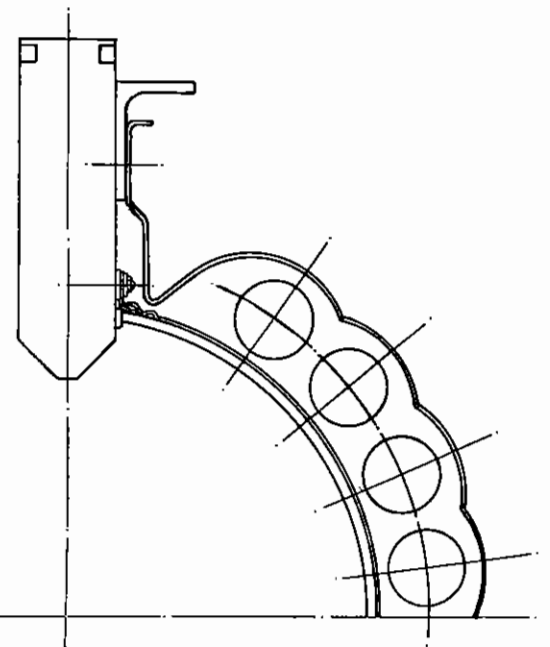


Fig. A6.3 Assembly details of Disc Amplifiers.

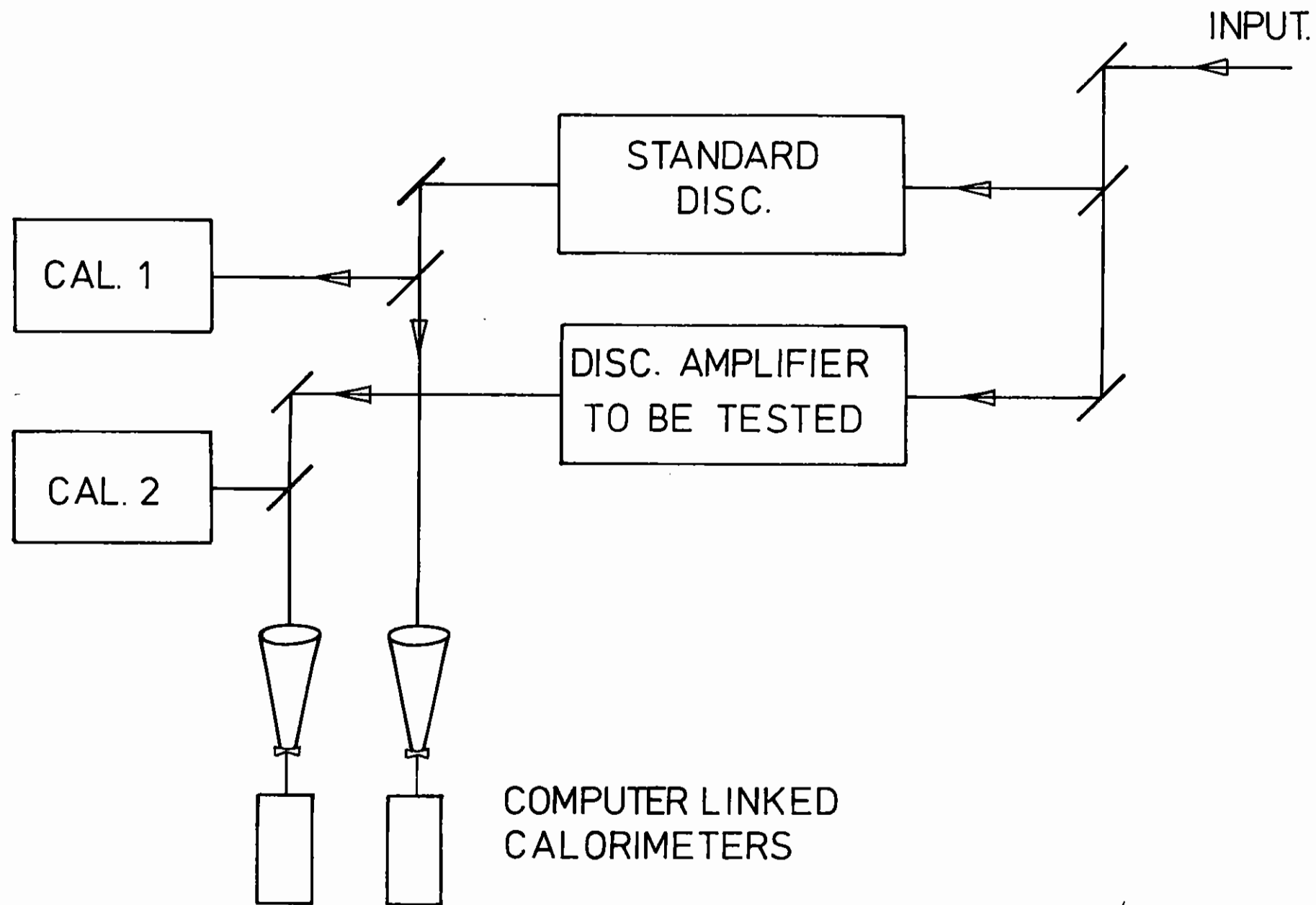


Fig. A6.4 Experimental layout for the disc amplifier testing programme.

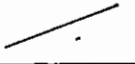
	Average Gain (± 0.2)	Average α (cm^{-1})	Reduction in α
20 lamps firing Normal chassis	8.5	7.93	
16 lamps firing (2 pairs of face pumped lamps removed)	5.7	6.45	19%
16 lamps firing (2 pairs of lamps removed next to chassis)	6.9	7.15	10%
20 lamps firing Modified Chassis	8.2	7.80	
16 lamps firing Old reflectors	7.2	7.30	6.5%
16 lamps firing New reflectors	7.45	7.44	5%

Table A6.1 Results of Test on $\phi 108\text{mm}$ Disc Amplifiers.

apertured by the chassis contributed only $1\frac{1}{2}$ -2% each.

When the modified chassis was tried with a twenty lamp amplifier no gain increase was seen. With the four chassis lamps removed a reduction in gain coefficient of 6.5% was measured. No observable increase in gain was measured when the modified reflector was installed. A summary of the results obtained is shown in Table A6.1.

In another test period the main measurements were repeated, the gain of the twenty lamp unmodified chassis was on this occasion compared directly with a sixteen lamp modified chassis. The average gains measured were 8.9 ($\alpha = 8.1\text{cm}^{-1}$) for the sixteen lamp amplifier. The reduction in gain coefficient was therefore 8%, while the reduction in gain was 16%. This result was in good agreement with the figures obtained in the earlier test.

It was concluded that although the sixteen lamp amplifier was more efficient (~5%) than the twenty lamp amplifier, it would still be unacceptable to use in any current or future system, as the loss of gain (16%) is too high.

A6.3.3 $\phi 150$ Disc Amplifier Development

J E Boon, P Gottfeldt and I Ross (RAL)

The prototype $\phi 150\text{mm}$ amplifier and power supply described in Section 1.3.6. of last years Annual Report have now been constructed. Figure A6.5 shows the preliminary mechanical assembly of the amplifier. The flash lamps have been manufactured and tested. These have a larger bore, 20mm, than the present 15mm bore tubes used in the $\phi 108\text{mm}$ amplifiers, and are intended to work at a design energy of 15.7KJ. This will be provided by a pulse forming network in the power supply using a capacitance of 130 μF charged to 22KV, driving 2 flash lamps in series through a 315 μH inductor. The design current peaks at 7.1KA after 260 μs . The anticipated lifetime to 70% of original light output is 2.2×10^5 shots. The lamps were tested using the Flashlamp Test Facility described in Section 1.3.7 of last years Annual Report. For these tests the energy per lamp was increased by 50% to 23.6KJ, being tested singly using a pulse forming network with 254 μF capacitance

charged to 13.6KV and driven through a 200 μ H inductor. This gave a peak current of 8.6KA after 290 μ s. All the lamps were subjected to 50 shots in the test facility. There were no failures.

A6.3.4 Laser diagnostic System Development

C Danson, B Espey, A Raven, D Smith and R Wyatt (RAL)

A6.3.4(a) Pulse Profile Measurements

A type 675 Hadland Photonics S1 streak camera has been purchased during the current reporting year, and has been tested in the laser area. Previously, short-pulse duration information was available only using an autocorrelator. Since this latter technique relied upon multiple pulses and requires prior knowledge of the pulse shape to yield an unambiguous result, the method is rather limited, particularly as many experiments now require a non-standard pulse-shape.

When fully interfaced, the new system will provide pulse shape information on a single shot basis, of both long and short pulse outputs at various points along the laser chain, with digital and graphical output available through the 4080 system computer via a CID camera and frame store combination.

Single mode optical fibres bring the pulses from various parts of the laser system to the slit plane of the streak camera. Careful cleaving of the fibres with a diamond cutting tool as well as careful positioning in the laser beam was found to be essential to produce good quality, repeatable images at the streak camera output. Further refinements in positioning are being considered to reduce the risk of pulse distortion and/or broadening in the fibre optic cable. Eventually about 10 fibres will be installed but at present only 3 are in use. Two bring pulses from the oscillators which can be viewed on a three second cycle for oscillator and switchout optimisation. A third fibre optic cable has been installed at the output of the system to examine the effects of gain saturation and etalon modifications to the pulse profile. Figures A6.6(a) and (b) compare a high and a low energy shot at

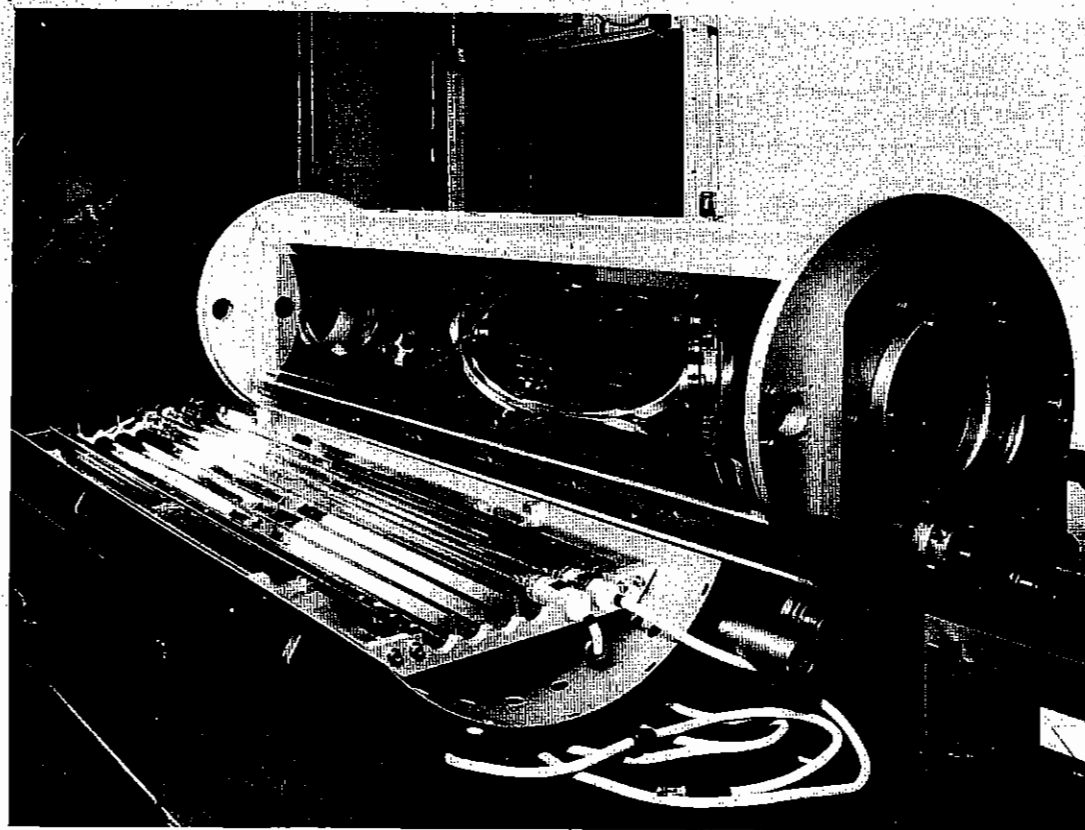


Fig. A6.5 Preliminary mechanical assembly of $\phi 150$ disc amplifier.

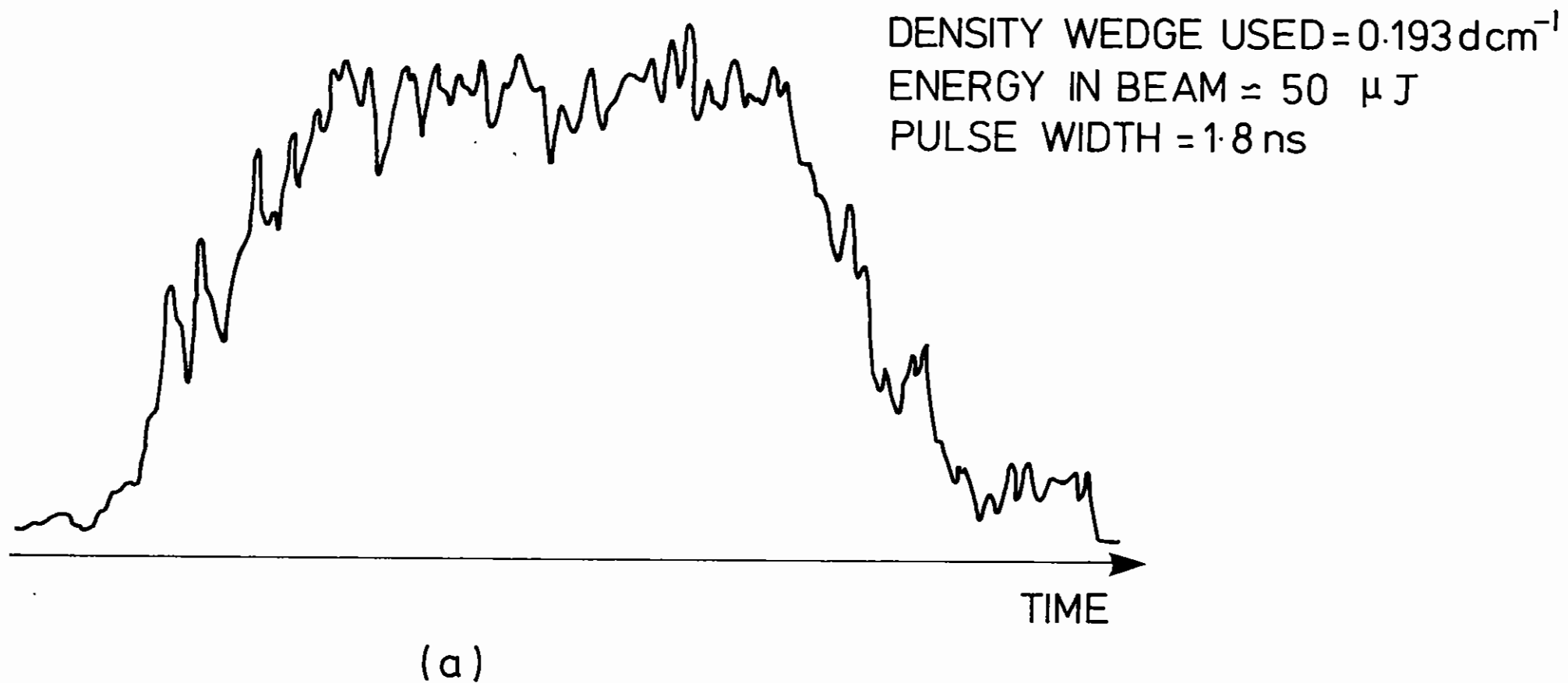
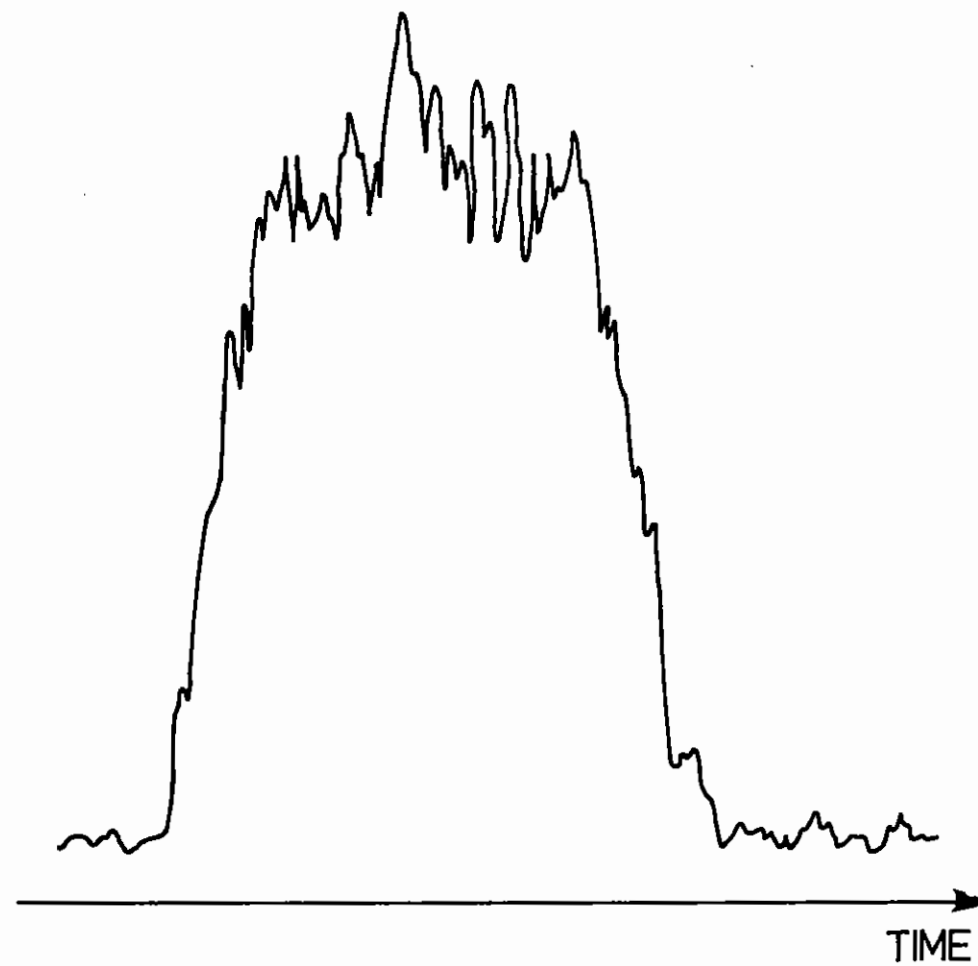


Fig. A6.6(a) Long Pulse Streak Record at Low Energy.

DENSITY WEDGE USED = 0.06 dcm^{-1}
ENERGY IN BEAM $\approx 80 \text{ J}$
PULSE WIDTH = 1.8 ns



(b)

Fig. A6.6(b) Long pulse streak record at relatively high energy.

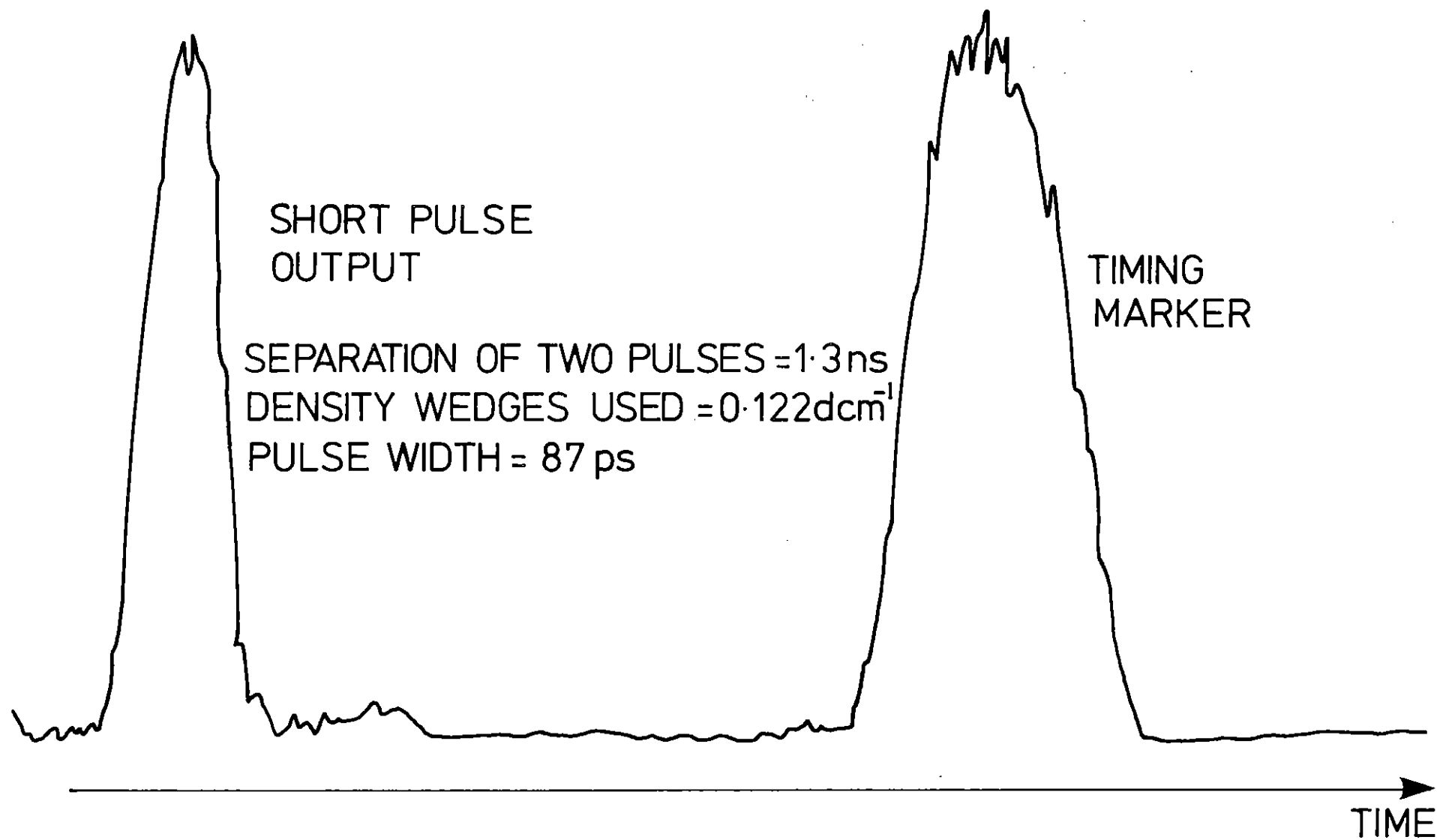


Fig. A6.7 Short pulse oscillator output.

the output of the system. At the 80J level no saturation is evident. A feature of the traces is the flatness of the pulse over the 1.8ns width, showing good switchout operation. Figure A6.7 shows a microdensitometer trace of the short pulse oscillator output, together with a timing mark provided by another fibre. The increased width of this pulse is due to the superposition of two short pulses. Figure A6.8 shows the pulsed shape that was required for the Carbon VI recombination X-ray laser experiment. It was produced using a prepulse etalon in front of the retro-mirror of the driver disc amplifier. All these traces were produced by using HP5 film in direct contact with the phosphor of the streak camera intensifier.

Experience to date with the fibre optic/streak camera combination indicates that this will provide invaluable information for both oscillator alignment purposes and target area diagnostics.

An Hakuto MV200 framestore with capacity to hold two camera frames was installed during the year. The framestore is triggered on a three second oscillator cycle to capture the streak camera image recorded by a GE 2506 CID detector. At present the image can be examined visually by means of a monitor in the control room but when the system is fully operational, data will be processed by the GEC 4080 computer. Once the data has been acquired by the computer, pulse heights, durations and rise times will become available immediately after a shot has been fired, and will be displayed along with the present laser status information. Development work on this system is at an advanced stage.

A6.3.4(b) Energy Monitoring Developments

Energy monitoring at various stages of the laser chain is important to measure amplifier performance and for fault diagnosis.

At the moment, there is no energy monitoring along the rod amplifier chain. The only positions available for energy monitors are behind mirrors and reflections from rod faces and the energies at these points vary from $\sim 250 \mu\text{J}$ to $\sim 100\text{mJ}$. The volume absorbing disc calorimeters used at present are not good enough to measure such low energies, so development

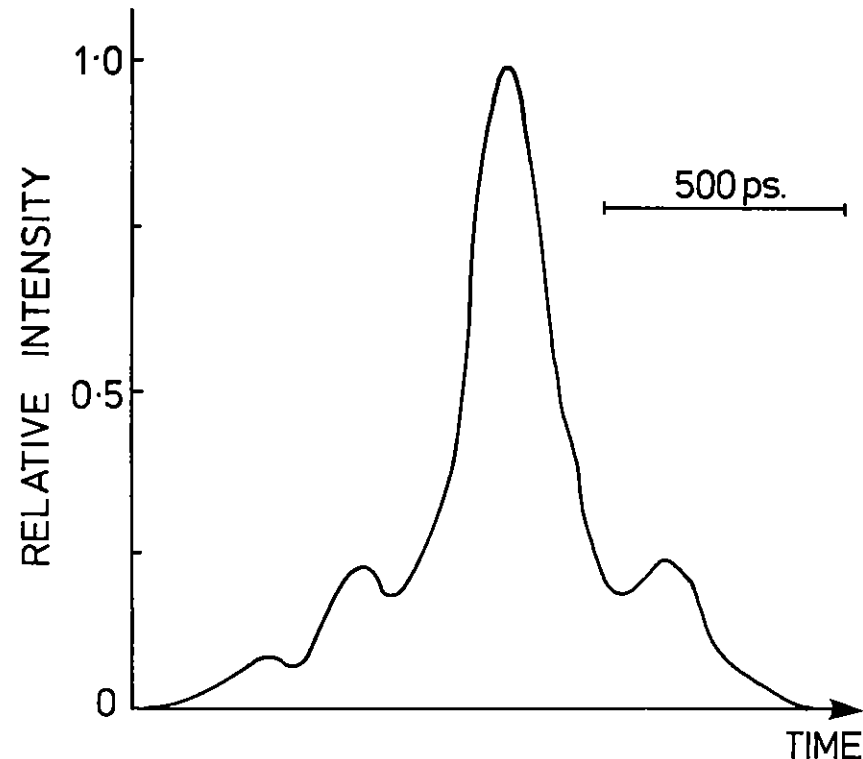


Fig. A6.8 Temporal pulse shape used for carbon VI experiment.

work on thermistor based detectors has been carried out.

Thermistors were supplied, already mounted in pairs on a T05 type transistor base by Thermometrics Inc. The basic scheme consists of connecting the two thermistors into a bridge circuit, with one exposed to the laser pulse and the other acting to compensate for changes in ambient temperature.

Using the bridge circuit in Figure A6.9 the output voltage from the bridge for a 1mJ laser pulse (using a YAG oscillator, similar to one of the glass laser oscillators) was of the order of 0.2mV. Due to this low output, very high precision amplifiers were needed. Various tests were carried out and it was decided, that to monitor energies of the order of 250µJ, a commercially available precision amplifier would be necessary, due to the high electrical noise in the laser area.

The amplifier chosen used was a Burr-Brown 3630. This amplifier is specifically designed for use on such bridge circuits and was found to be ideal. Figure A6.10 shows the final circuit used. A sample-hold circuit was needed because the computer reads three seconds after the laser fires and the thermistors response was much faster than this (see Figure A6.11).

The calorimeter was mounted in an aluminium cylinder and placed in the laser area to monitor the long pulse oscillator output. It was tested for linearity against a calorimeter at the end of the rod amplifier chain (see Figure A6.12) and finally calibrated against a Rofin energy monitor.

The calorimeter is now connected permanently to the computer and working successfully. Further development on precision amplifiers has been successfully carried out, and four more calorimeters to go at various points down the rod amplifier chain, are being manufactured. The amplifiers are of a similar design to that of the Burr Brown 3630 and the op-amps used in construction are FET input LF356's. This type is used for their very good Common Mode Rejection (> 100dB).

A6.3.4(c) Short pulse oscillator Prelase Monitor.

Monitoring the prelase waveform from an actively mode locked Q-switched oscillator gives essential information about the quality of alignment and cavity length timing. In order to obtain a permanent record of this waveform a Phillips PM3305 storage oscilloscope was acquired during the year for use as a prelase monitor. This scope has a 4k internal memory which can hold both a digitized version of the displayed waveform and a list of the front panel settings. By using an interface option (PM8955), the scope has been connected to an IEEE 488 interface bus (GPIB) linked to the GEC 4080 computer. Accessing the stored information begins after the firing of a main disc shot and the data points, on receipt by the computer, are formed into running means of the maxima and minima of the waveform so that the prelase envelope can be obtained. The envelope is then plotted along with the transient digitizer data of the short and long pulse oscillator output and also stored on disk for future reference. Figure A6.13 shows such output. These data provide both a record of the laser's performance and an indication of when oscillator adjustment is necessary.

A6.3.5 Rod amplifier flash lamp test facility

R Wyatt (RAL)

The facility for testing the large flashlamps used in the disc amplifiers has been invaluable and problems caused by installation of faulty lamps have now been effectively eliminated. Based on this experience it was decided to construct a test facility for the smaller lamps used in the rod amplifiers in order to eliminate faulty lamps and identify those with low output and triggering problems.

The lamp test housing was purpose built to accommodate the two lengths of lamp used in our system; 20cm and 10.5 cm. In practice both are fired from a 480 uF source, through either a single long lamp or two short lamps in series. To enable single short lamp testing a Quantel Capacitor Unit of 240µF was used, and a resistance wire of finite value used to replace the copper wire connection from lamp to ground to maintain critical damping.

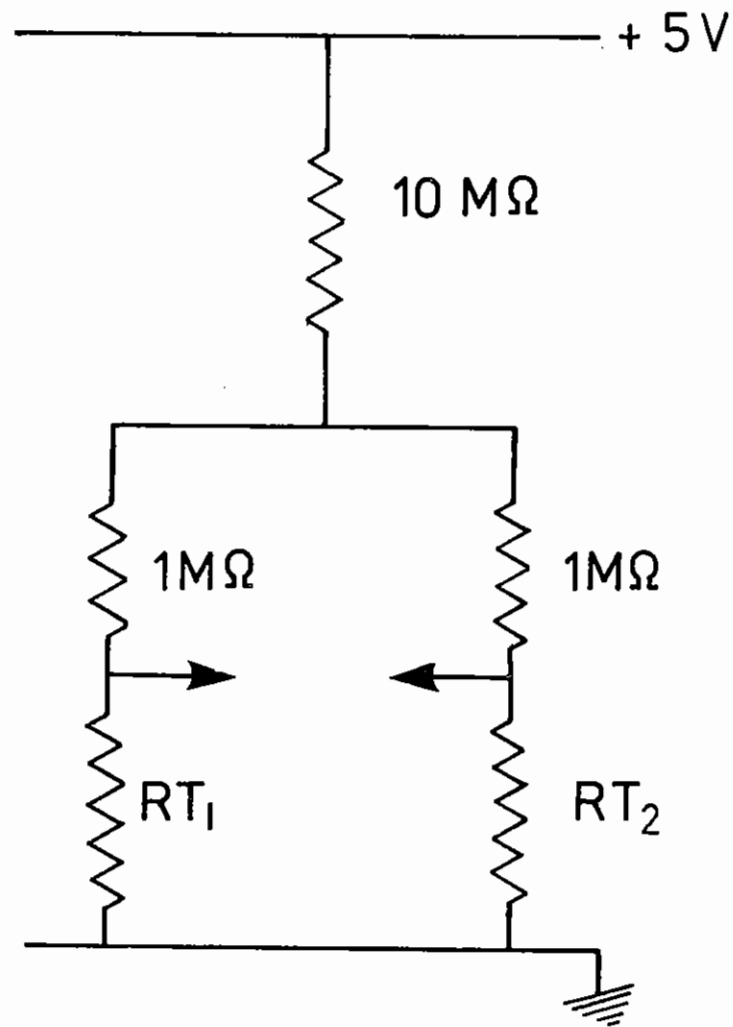


Fig. A6.9 Simple bridge circuit used for thermistors during initial energy monitor tests.

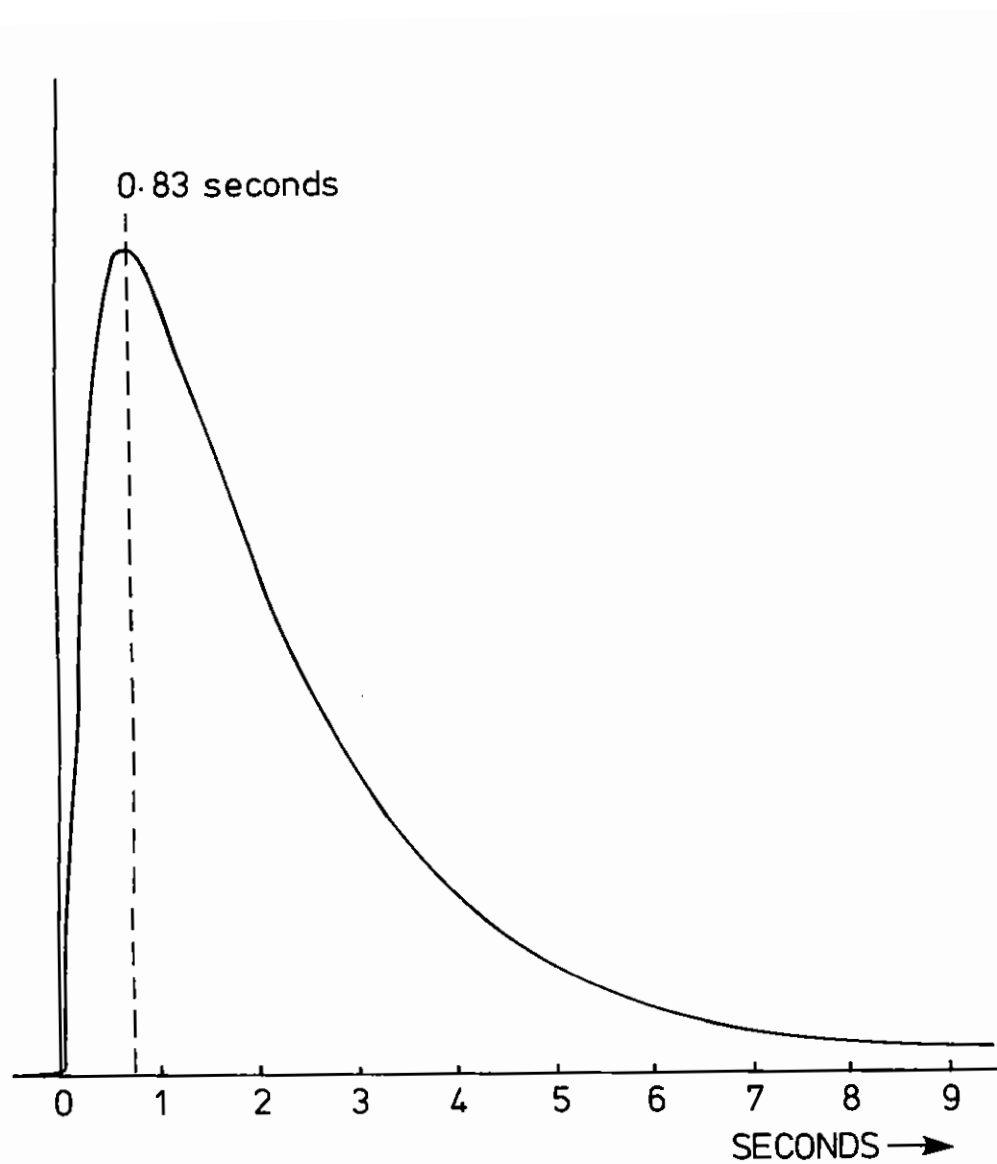


Fig. A6.11 Response of thermistor for a 1mJ input laser pulse.

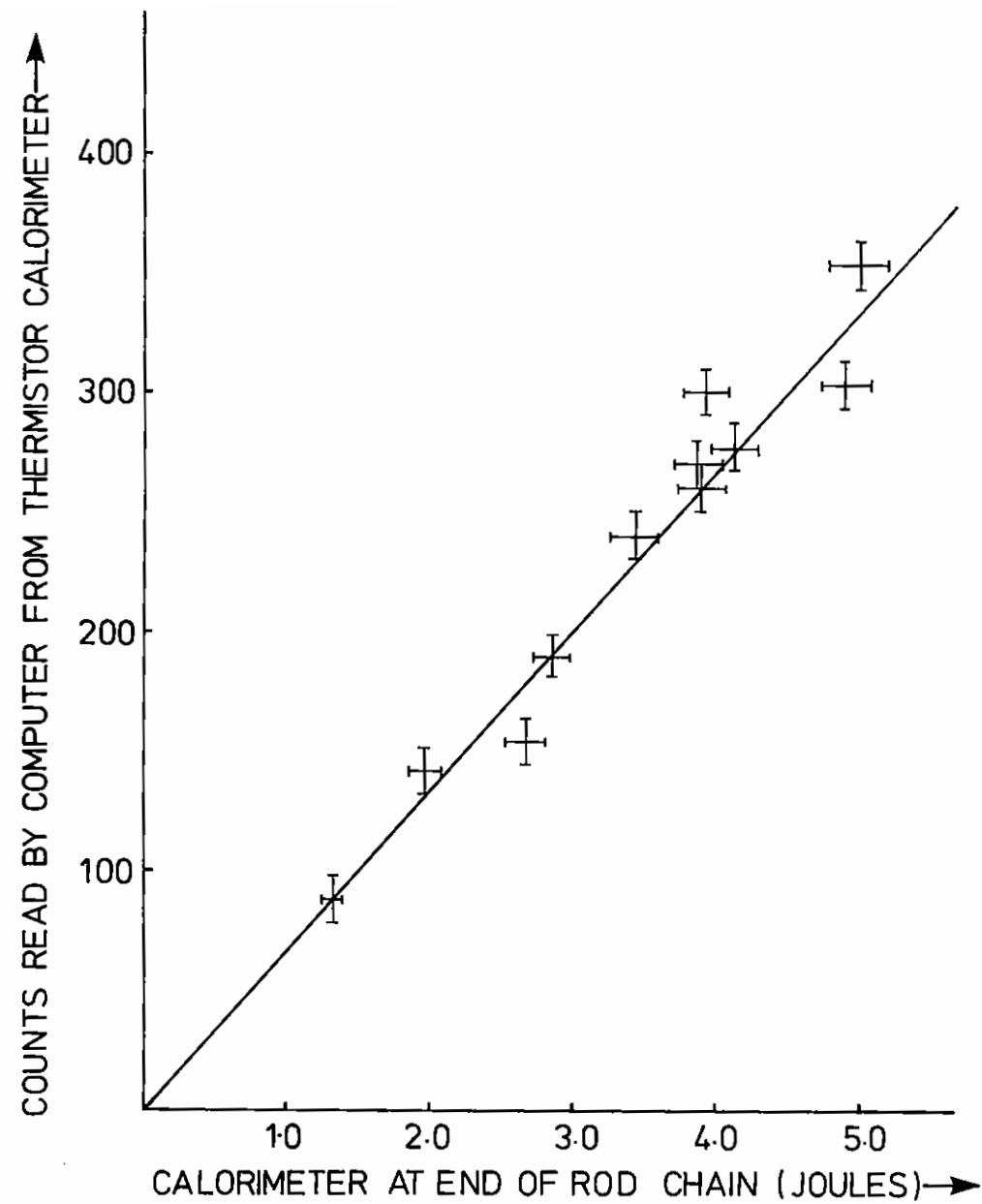


Fig. A6.12 Results of thermistor calorimeter linearity tests.

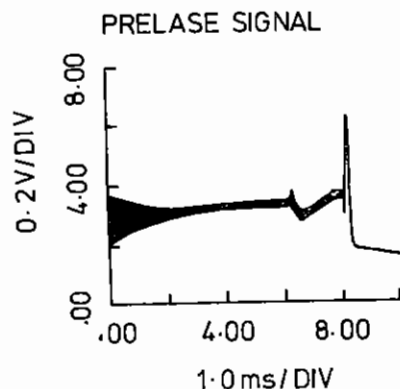


Fig. A6.13 Output from short pulse oscillator prelease monitor.

Figure A6.14(a) shows a schematic diagram of the system. Current through the lamp circuit was monitored using a Rogowski coil while the useful output of the lamp was recorded using a fluorescence monitor.

To provide the $1\mu\text{m}$ fluorescence signal a disused 9mm amplifier rod was built into the unit. The fluorescence from the rod is coupled to a PIN diode through a fibre optic link, and a $1\mu\text{m}$ filter. A second fibre optic link directs flash light to a PIN diode in a lamp failure detection circuit. The Current and fluorescence standards were determined from tests on a number of new lamps, carried out in the test facility.

Control and power is provided from two Quantel units, CU 401 & PU 420, which will run on a two minute rep rate at pre-selected energy levels, within the normal operating range of our system. In event of lamp failure the system has been designed to dump and shut down automatically.

The test facility also allows variations in the lamp trigger signal to be carried out in order to identify lamps which may be developing triggering problems.

Output waveforms are recorded on a storage scope and compared with known standards. Figure A6.14(b) shows typical output current, total light and fluorescence waveforms.

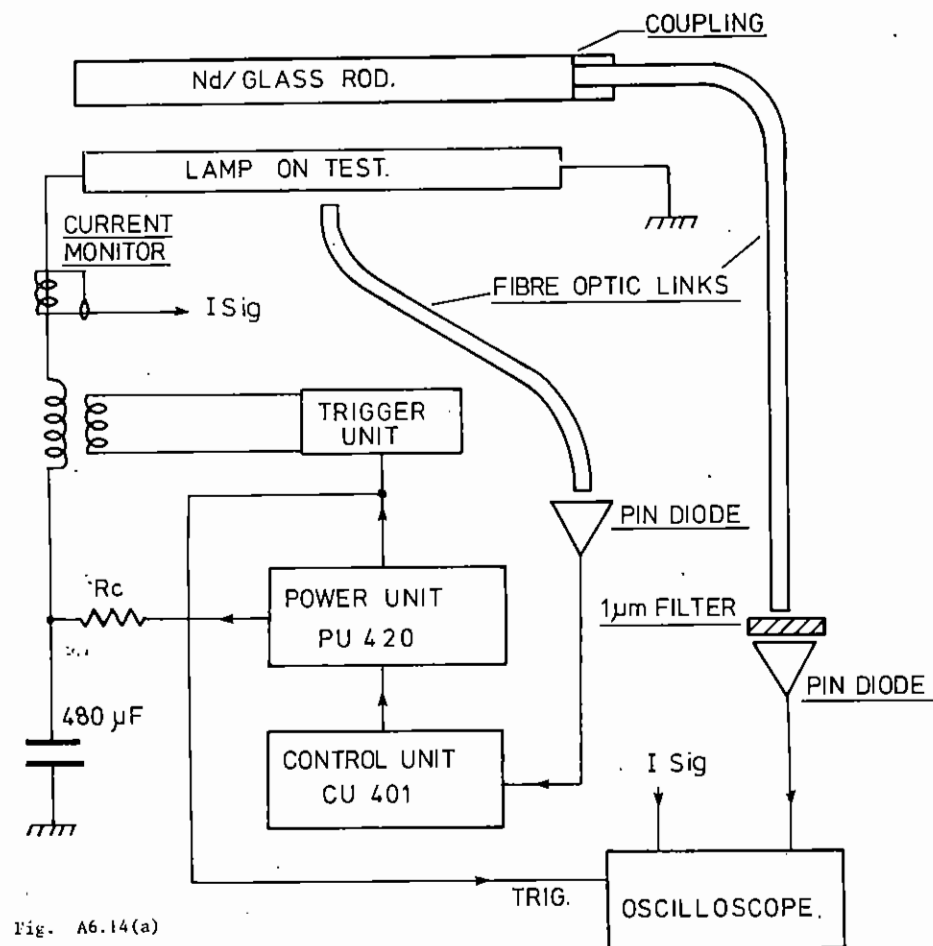
Many tests have been performed on both old and new lamps. These have led to significant improvements in rod amplifier operating reliability.

A6.3.6 Damage Test Facility

N Allen and C Danson (RAL)

A simple facility for determining the damage threshold of the various coatings and materials used on the main glass laser system has been set up. This facility is positioned after the vacuum spatial filter following the final rod amplifier of the outer beam line and can be accessed using a remotely controlled sliding mirror. The set up consists basically of a long

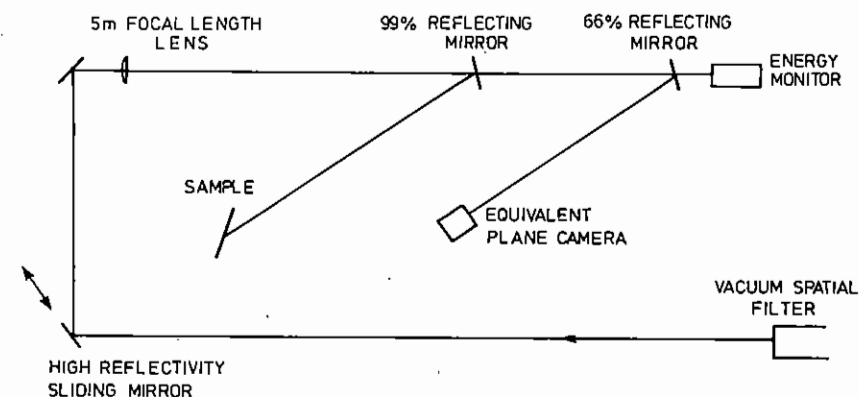
SCHEMATIC OF FLASH LAMP TEST FACILITY.



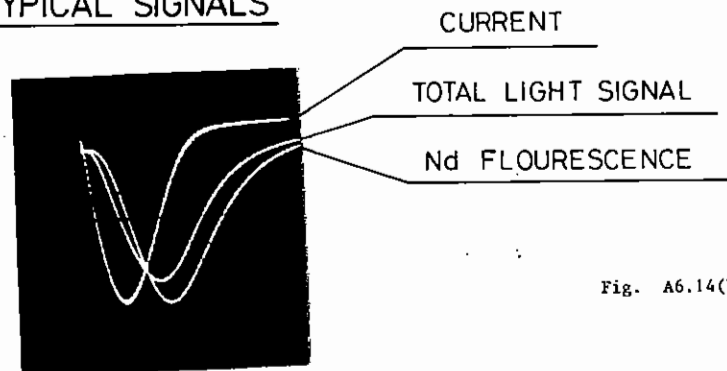
SAMPLE	'P' POLARISATION		'S' POLARISATION	
	SINGLE SHOT	MULTIPLE SHOT	SINGLE SHOT	MULTIPLE SHOT
CRYOLITE	> 7.7 > 9.1	~ 7.0		
QUEENS UNIV BELFAST ^⑤	~ 3.0			
④	~ 6.0			
SCHOTT NEUT. SOL'N.	> 12.5	~ 16		
TECH. OPT (T.O.P.4)		> 8.5 < 12	< 2.9	
O.C.L.I. POL'R. S/44	~ 5.3			
SAMPLE 148	5.5			> 20.5

SAMPLES WERE PLACED AT 56.5° ANGLE TO THE NORMAL TO THE BEAM

Table A6.2



TYPICAL SIGNALS



focal length lens which focusses the beam onto the sample under test. Limitations of space necessitated the use of a folding mirror as indicated in Figure A6.15. The small amount of transmission through the folding mirror was utilised to measure the beam energy and obtain equivalent plane photographs. The facility was calibrated and some initial tests were carried out on thin film polariser coatings, single layer cryolite A/R coatings and a sample of Schott neutral solution treated BK7 glass.

The criteria for onset of damage was basically a visible increase in scattered light, when viewed through a 10X microscope. Initial results are as indicated in Table A6.2.

A6.3.7 Computer control for driver system

P Gottfeldt and D Pugh (RAL)

The touch screen concept for laser control was briefly described last year. In order to provide touch screen control for individual amplifiers of the Quantel/J.K. Laser driver system certain modifications, particularly to the Quantel system, were required. Furthermore, the Quantel system was never designed with a true dump facility making it impracticable to abort a laser shot part way through a charge sequence. In view of our efforts to allow laser users to control their own shots it was felt desirable to provide a genuine dump facility.

ON/OFF switching of individual Quantel amplifiers and Faraday Isolators required the design of a special solid state switching circuit for the charging supply control units. The Quantel system operates from a 380V mains supply and it was found that standard solid state relays for this voltage were unsuitable for the low power circuits involved due to excessive leakage current in the off state. The circuit finally devised uses 2 SCR's back to back triggered from a UJT which provides gate pulses via an isolating transformer. The UJT is fed from a D C supply which is also the switching signal.

The trigger pulse for the Quantel amplifiers is superimposed on the 24V

supply which is also used to hold in a dump relay. Simple removal of the 24V to provide a dump for the Quantel system generally causes the amplifier flash lamps to fire. To prevent this from happening it was necessary to interpose another relay contact in the dump relay circuit allowing the dump relay to be released without removing the 24V supply. The charging supply control units required additional modifications to prevent them from issuing trigger pulses to the charging supply triacs after the charge command had been removed. The dump sequence now resets the universal charge command latch which then removes the charge signal from all the charging supply control units and stops trigger pulses reaching the charging supply triacs. Simultaneously the new control relays for the dump system are de-energised causing the stored energy to be dumped.

A CAMAC - driver amplifier unit has been constructed to give the GEC 4080 control computer access to the individual amplifiers and Faraday isolators, and to provide an input for the dump signal.

A6.3.8 Enhancement of Computer control and data acquisition

P Gottfeldt and C Reason (RAL)

The most significant enhancement of the control system has been the addition of a second CAMAC branch to the GEC 4080 computer control system. The branch will be used to provide a data link between the GEC 4080 and the new target area complex. It will also allow for control of the laser from the south control room.

The existing parallel branch uses a 132 way cable and despite the use of differential branch extenders it appears to be susceptible to occasional electro-magnetic interference. Open circuits at the connectors have also occurred from time to time and it was considered unwise to extend this data highway by some 40m. Instead a second branch was added to the system built to the CAMAC serial standard EUR 6100e running in bit serial mode and using U-port adaptors for fibre-optic transmission.

The U-port adaptors, manufactured by CAMTEC to a JET design, are housed in

19" rack mounting units and incorporate battery back-up in case of a local mains failure, by-pass and loop collapse mechanisms and automatic transfer to a standby fibre link. CAMAC crate control in TAW is provided by an L2 crate controller which permits the use of an auxiliary controller in the same crate. This arrangement is used to provide an autonomous control system for target and beam alignment in TAW with access to data from the GEC 4080. The serial branch driver and serial branch handler reside in the GEC systems crate together with the parallel branch driver. In this way it has been possible to make the two branches compatible although significant changes to the existing software were necessary. The serial branch will be used for most new applications but it is proposed to retain the parallel branch for applications where the high data transfer rate it offers may be required.

During the year, the Electric program on the IBM mainframe was withdrawn and the link between the GEC 4080 and the IBM was changed from HASP workstation to VNET. This gives an easy path between the 4080 and a CMS virtual machine but has entailed changes in the way files are exchanged and jobs submitted. This has also resulted in a loss of terminal access to the IBM from the control room via the 4080. It is hoped to replace this with a separate link via a JNT PAD in the future.

A6.4 Target Preparation and Development

C Brown, B Child, M Davies, N Prior, P Rumsby and S Whittaker (RAL)

As well as producing close to 2000 hollow shell and planar targets for experiments the target fabrication group has been involved in a wide variety of research and development work during the past year.

A major improvement has taken place in the area of glass shell selection. Up to now all hollow glass shell targets used for experiments at the laboratory have originated from raw material supplied by 3Ms company. Quality of this material is very low and considerable time has been devoted to the selection of suitably uniform shells. In addition, wall thicknesses

above about 2 μ m were not available in this material and consequently targets for ablative compression experiments required additional layers of material to be deposited to produce final targets with adequate thickness to minimize preheat. Both these limitations have now been overcome as 20 batches of high quality shells of both high and low aspect ratio have been ordered and received from KMS Inc. Manufacturing of these shells take place in drop towers in which the process stages are closely controlled leading to shells of extremely high uniformity. About 50% of the shells in any batch are of target quality compared to the much less than 1% for the previously used material. Batches have been obtained both DT filled and empty with wall thickness ranging from 0.7 to 6 μ m. Time spent on target selection has now been dramatically reduced.

Coating techniques using metallic slit masks have been developed in order to produce novel planar targets for two different experiments. For 2 beam colliding shock wave experiments two 15 μ m wide by 300nm thick Bismuth marker stripes have been deposited on opposite sides of a 30-40 μ m thick polystyrene foil which has subsequently been coated on both sides with 10 μ m of N type parylene. The exact registration of the two stripes on opposite sides of the foil was critical to the success of the target design and this was achieved by optically aligning two masks with slits 15 μ m wide by 6mm long and then placing the polystyrene substrate between for coating. For transport experiments 2 μ m thick aluminium microdot targets have been deposited on polymer substrates by using two crossed slit masks. Targets with dots 15 μ m and 30 μ m square complete with marker lines have been prepared by this method as shown in Figure A6.16.

Ultra violet photo ablative etching techniques have been pursued further. An ArF laser ($\lambda = 193$ nm) has been acquired and work has continued on both mask and interference techniques for etching polymer targets. 20 μ m periodicity patterns have been etched on polystyrene shells for two series of Rayleigh-Taylor experiments. Etch depths from 0.5 to 2 μ m have been produced with a high level of reliability and reproducibility. Figure A6.17 shows an example of such a target with a 2 μ m etch depth. About 100 such targets have been prepared by this method. For transmission X - radiography

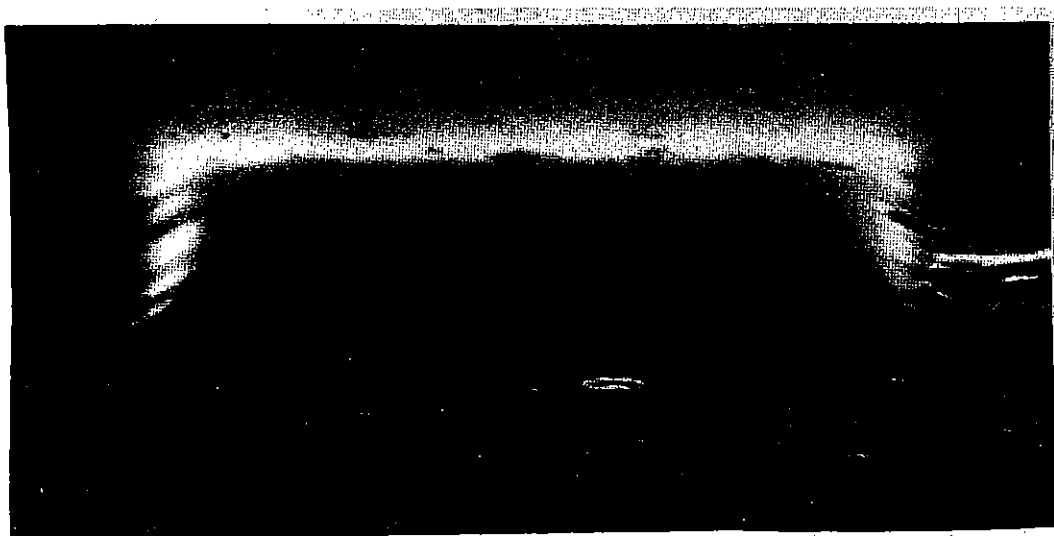


Fig. A6.16 Aluminium microdot target $15\mu\text{m} \times 30\mu\text{m} \times 2\mu\text{m}$ thick deposited on a polymer substrate.

experiments the rear of striped polymer shells have been removed or thinned (to reduce the X ray opacity) by irradiating them through slits or apertures. Using a line narrowed KrF excimer laser ripple patterns of $15\mu\text{m}$ periodicity and $0.5\mu\text{m}$ depth have been generated on planar Mylar targets using interference methods. This technique is being actively researched as it is significantly better than contact methods as sinusoidal ripples are produced and mask edge diffraction effects are eliminated. The photo ablative etching of low density polyacetylene foams has been investigated and found to be a very successful method for machining such material. In addition it has been found that thin parylene foils ($1\mu\text{m}$ or more) can be readily machined by this method to leave free standing structures.

Two developments have taken place in the manufacture of targets for X-ray laser experiments. For a recombination laser scheme carbon fibres have been prepared with diameters down to $2\mu\text{m}$ by reducing commercial $8\mu\text{m}$ fibres in diameter using reactive r.f. plasma etching techniques. For a resonantly photo-pumped neon laser scheme $3\mu\text{m}$ thick aluminium foils have been doped with neon by up to 10% by ion implantation at 80KeV.

A6.5 New Target Area Facilities

A6.5.1 Introduction

J E Boon, M H Key and W T Toner (RAL)

The building extension mentioned in last year's report has been occupied since the autumn and new equipment has been installed throughout the Winter. The layout is shown in Figure A6.18. Target Area East has been in use as an extension to Target Area 2 since the Autumn for complex experiments needing additional floor space, such as the Z-pinch experiment (Section A1.6). The versatile multi-beam irradiation facility planned for this area was delayed for financial reasons but detailed design work is now in progress. The 12-beam symmetrical irradiation facility which is being assembled in Target Area West is described in detail below. Since experience has shown the design of the present six-beam facility to be satisfactory the new chamber

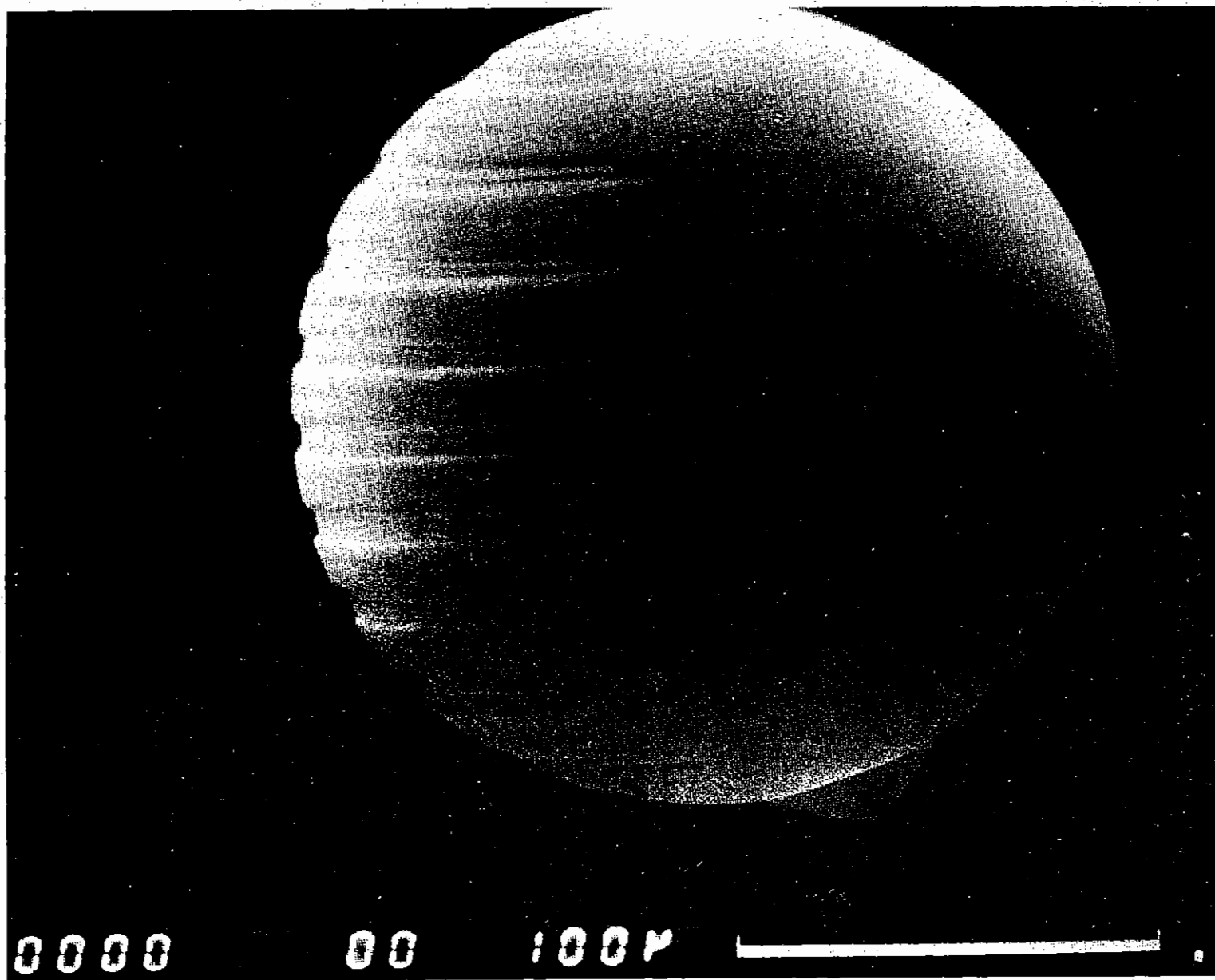


Fig. A6.17 Polystyrene hollow shell target etched $2\mu\text{m}$ deep.

and its equipment is similar to the old, but numerous improvements of detail have been made. The major departure from current practice is the incorporation of a microcomputer in the control system to simplify the operators job, help avoid errors and speed up the alignment process. Views of the facility during assembly are shown in Figures A6.19(a) and (b). The first experiment is scheduled to begin in June, six weeks after the last experiment in TAl. From then on target area staff and users will operate experiments in TAW and TA2 from the new South Control Room, linked via CAMAC to the GEC 4080. Two small dark rooms have been provided nearby at mezzanine level.

A6.5.2 Optical Arrangement of 12 Beam System

A R Damerell, C J Hooker and W T Toner (RAL)

Targets will be irradiated by twelve green or infra-red beams arranged in dodecahedral geometry and focussed by f:2.5 aspheric doublet lenses (Section A6.6) mounted in improved versions of the lens carriages used in the six beam chamber. A typical beam is shown schematically in Figure A6.20. The components of the system are controlled by the operator through the microcomputer system described later. The disc amplifier output is brought to TAW by a chain of mirrors the last two of which are steerable.

Local CW alignment beams are injected in TAW by means of slide mirrors. An achromatic television comparator station (one of six) provides near-and far-field images so that the operator can make the local beams colinear with the main laser alignment system. Harmonic crystals are placed just before the splitters which make twelve beams from six. The beams are then directed to the target chamber. A second television station (one of twelve) views light reflected from the target in near and far field. Light transmitted past the target is viewed by the television station in the opposite beam line. All beams have an absorbing glass calorimeter which collects light reflected from a tilted, uncoated, chamber entrance window. An equivalent plane monitoring system will be installed progressively during 1984. Image relay from the disc amplifiers to TAW is under consideration.

A6.5.3 12 Beam Target Chamber and Diagnostics

A R Damerell, C J Hooker and W T Toner (RAL)

The chamber is manufactured from two hemispherical mild steel dishes welded together and electroless nickel plated. It is slightly larger (1093mm internal diameter) than the six beam chamber. Improved versions of the present airlock loaded target positioning devices will be used. An orthogonal pair of fixed reference microscopes will enable targets to be positioned precisely and reproducibly. The vacuum system incorporates a Rootes type booster roughing pump for rapid pump down as in the TA2 system. In order to provide optimum access to the chamber for experimenters all pumps are installed below floor level and liquid nitrogen is fed to the diffusion pump cold trap through a vacuum insulated transfer line in floor ducting from a dewar emplacement at the side of the room.

The twelve beam geometry at f:2.5 offers forty percent increase in the solid angle available for diagnostics over the six beam f:1 system. Sixteen of the twenty possible maximum access directions between the laser beams are provided with full aperture ports and there are numerous smaller ports. A brief note specifying details of diagnostic access has been circulated to users. Further copies are available. All the present diagnostic systems will be recommissioned on the new chamber during the experiments scheduled for the summer of 1984. Backlighting and optical probe beams will be provided as at present. The increase in diagnostic access will permit more diagnostics to be run simultaneously than is the case now and should also make it easier to align backlighting targets.

A mark II version of the inertial positioning system for stalkless targets described in last year's report is being developed in collaboration with MPIQ Garching and Queens University Belfast for use in the autumn of 1984.

When the large aperture disc amplifier development is complete an enhanced backlighting capability will be provided. Further development of the optical probe beam system is under consideration. Upgraded diagnostic instruments will be purchased as funds permit.

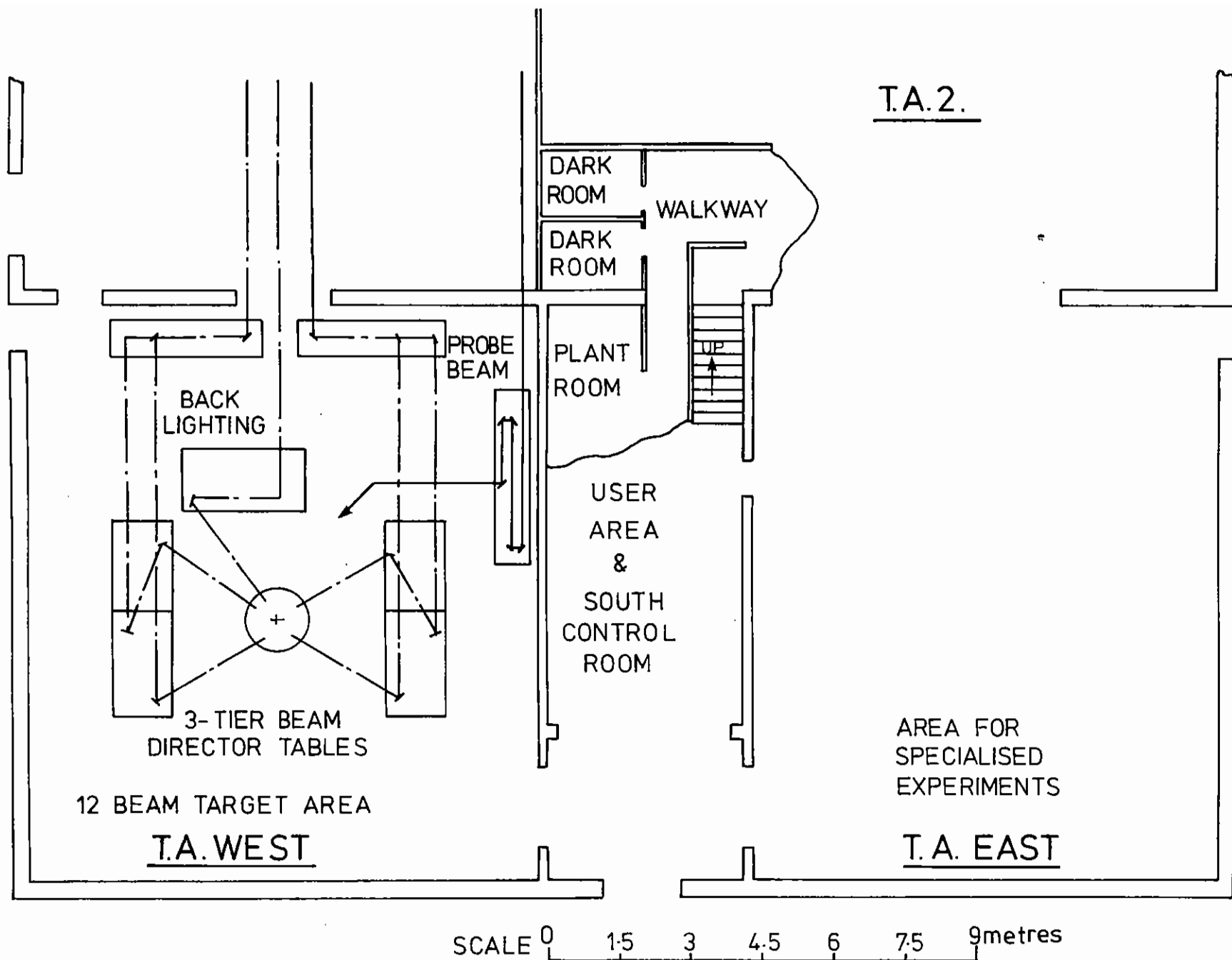


Fig. A6.18

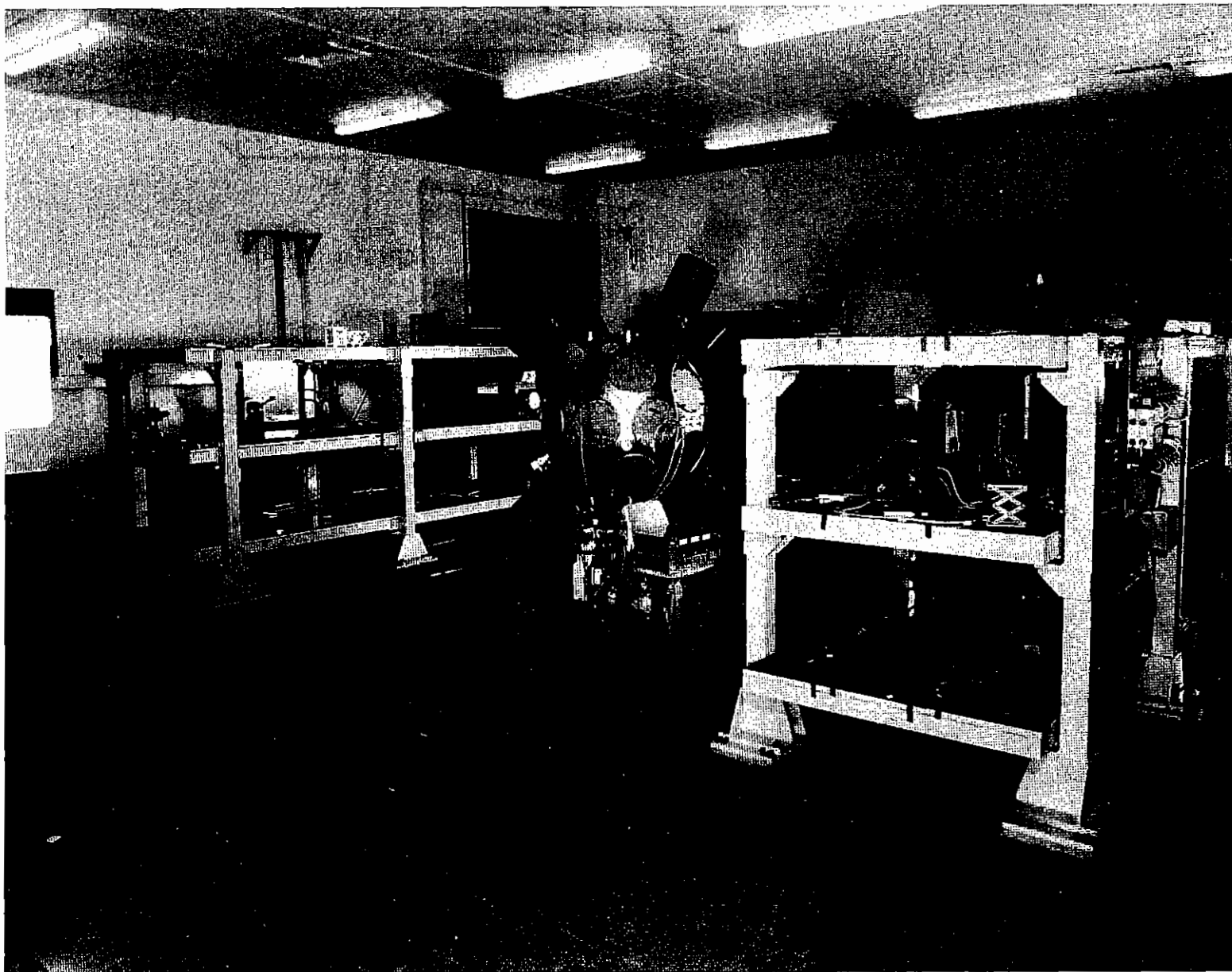


Fig. A6.19(a) View of Target Area West during assembly.

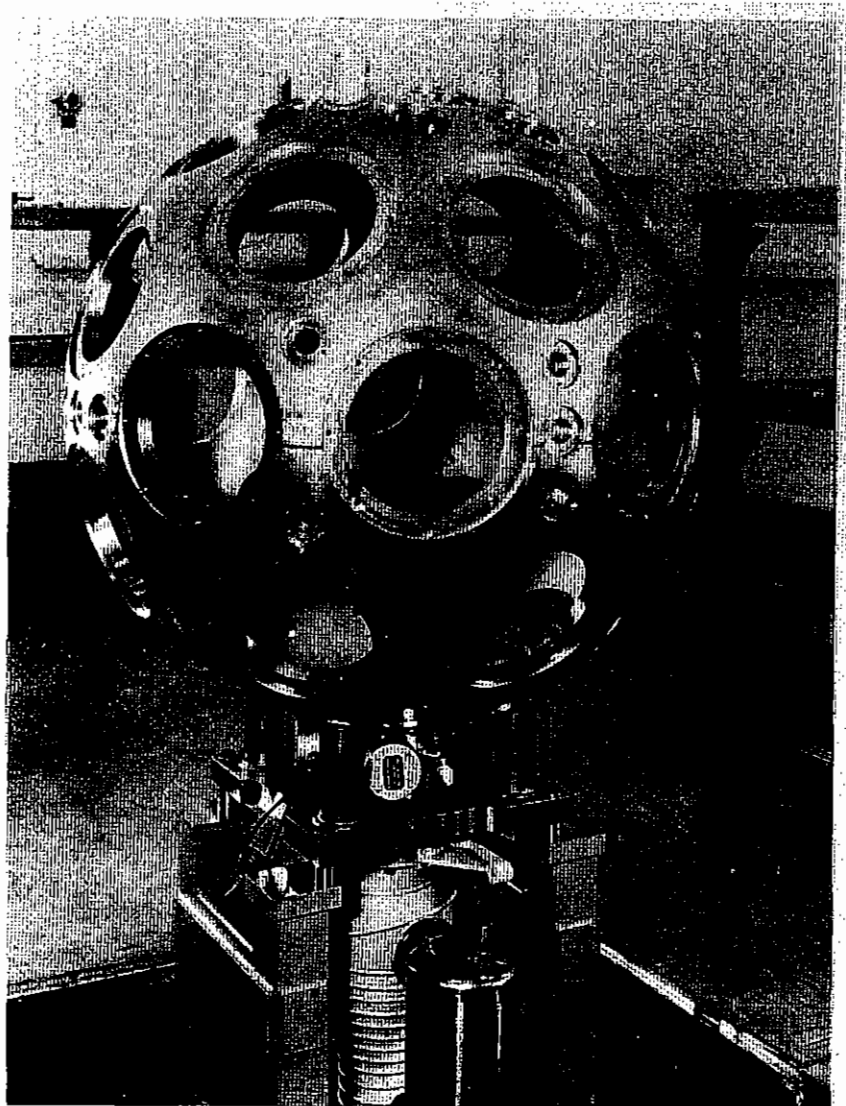


Fig. A6.19(b) 12 beam target vessel before assembly of components.

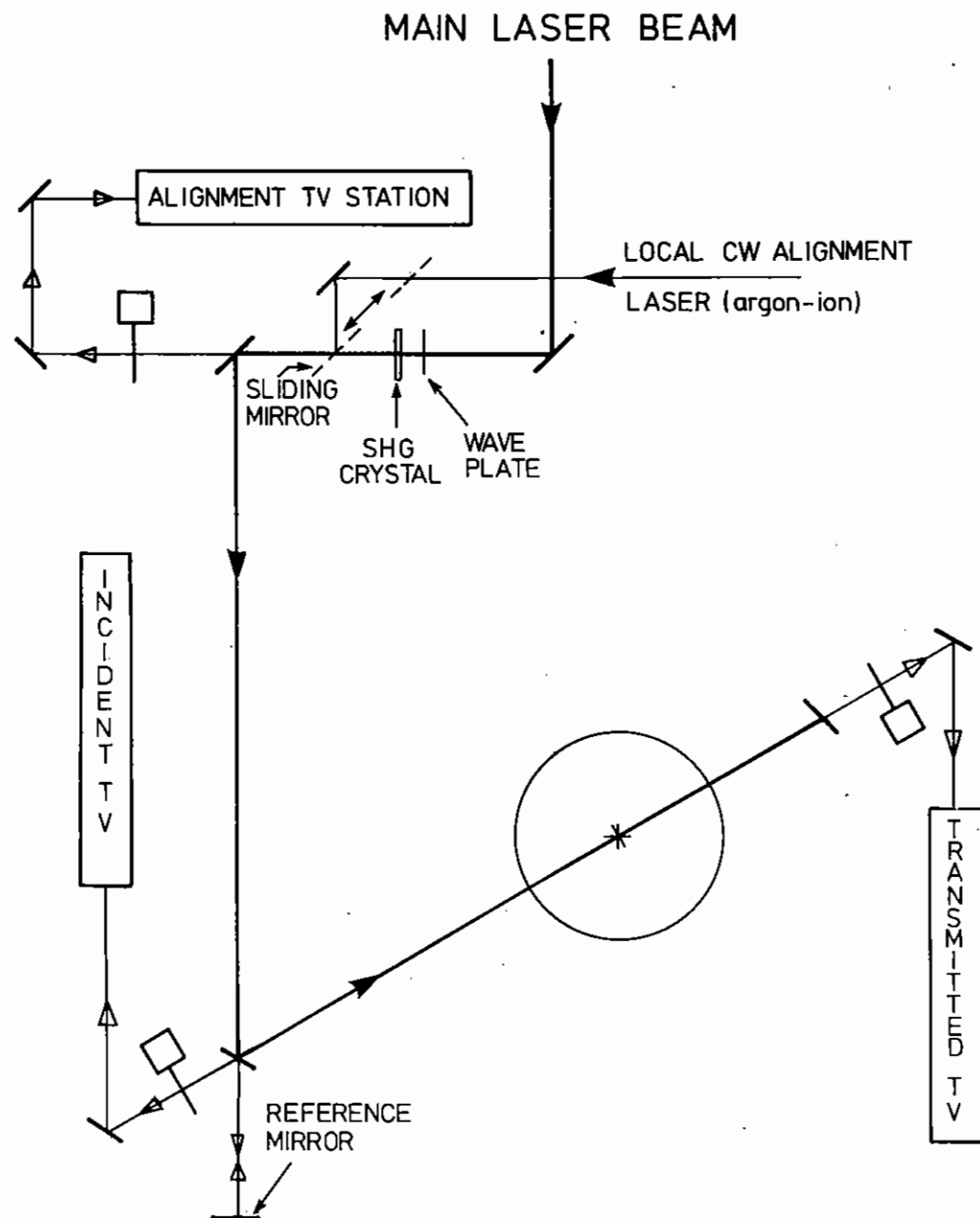


Fig. A6.20 Schematic of typical beam line of the 12 beam system.

In the long term it is planned to install further disc amplifiers in the laser so that all twelve beams may be driven directly at higher energy.

A6.5.4 Control System: Introduction

W T Toner (RAL)

The target area operator has the best part of two hundred remote controls of ten types at his command with status information coming from each and television signals coming from twenty sources. Although all beamlines are similar at the hardware level, they are not identical. Alignment procedures are basically similar for most experiments but again not identical. A wholly manual control system of the type used in Target Area 1 would inevitably lead to confusion, errors and damage to equipment. The CAMAC based microcomputer control system described here has been designed to simplify the operator's job as much as possible while preserving the flexibility to cope with varied experiments. It also permits a modular approach to the hardware which should make it possible to modify and extend the system in the light of experience. All the complexity resides in the software. This is written so that the operator is guided through the standard alignment sequence by a menu-style programme but has the freedom to depart from it and control individual devices in any order if he chooses.

A6.5.5 Control Hardware

M Forster, P Gottfeldt, W Lester and C J Reason (RAL)

The target positioning devices, the vacuum system and all adjustable optical components, shutters and television stations are controlled through the system shown schematically in Figure A6.21. At its heart is a DEC 11/21 Falcon microcomputer housed in CAMAC as an "Auxiliary Controller" and connected to the target area equipment via CAMAC modules and a custom built

modular interface system. Communication with the operator is through a VDU and keyboard. The Falcon can communicate with the GEC 4080 computer used for the main laser system via a fibre optic serial branch. The CAMAC crate is also used under 4080 control to drive a terminal and plotter in the South Control Room as well as to house data acquisition modules controlled and read by the 4080.

Communication with the 4080 is bit serial (to CAMAC standard EUR 6100e) and implemented in fibre optics using a master and slave U port adapter made to the JET design. This method was chosen to avoid noise pickup and other problems experienced in the parallel branches elsewhere in the CLF system. L2 rather than L1 designation of the crate controller allows a second "Auxiliary Controller" to be installed in the crate, connected to the L2 controller by the "Auxiliary Controller Bus" (ACB), through which access can be gained to the CAMAC modules in the crate.

The L2 controller acts as arbitrator between the two sources of commands using the "Auxiliary Controller Lockout" (ACL) signal as defined in EUR 6500eN.

The device control system is based on 3U high DIN 41494 specification sub-racks configured to hold up to 21 cards. The system can accommodate up to 15 racks. The cards in them contain the drive electronics for the devices, which are connected through the front panel. This gives great flexibility in allowing cards controlling different types of device to occupy the same rack. The cards receive information from the backplane which comprises 4 write lines, 4 read lines and shot select, clock and power lines. The action to be performed is set up on the write lines, the appropriate shot is selected and the action is performed on receipt of a clock pulse. The readlines are used to send limit and position information back to the computer. A buffer card in the right hand slot in each rack buffers and stores information coming from the computer and also decodes the 8-bit word containing the shot select information. As a result only one slot at a time can be selected in a rack.

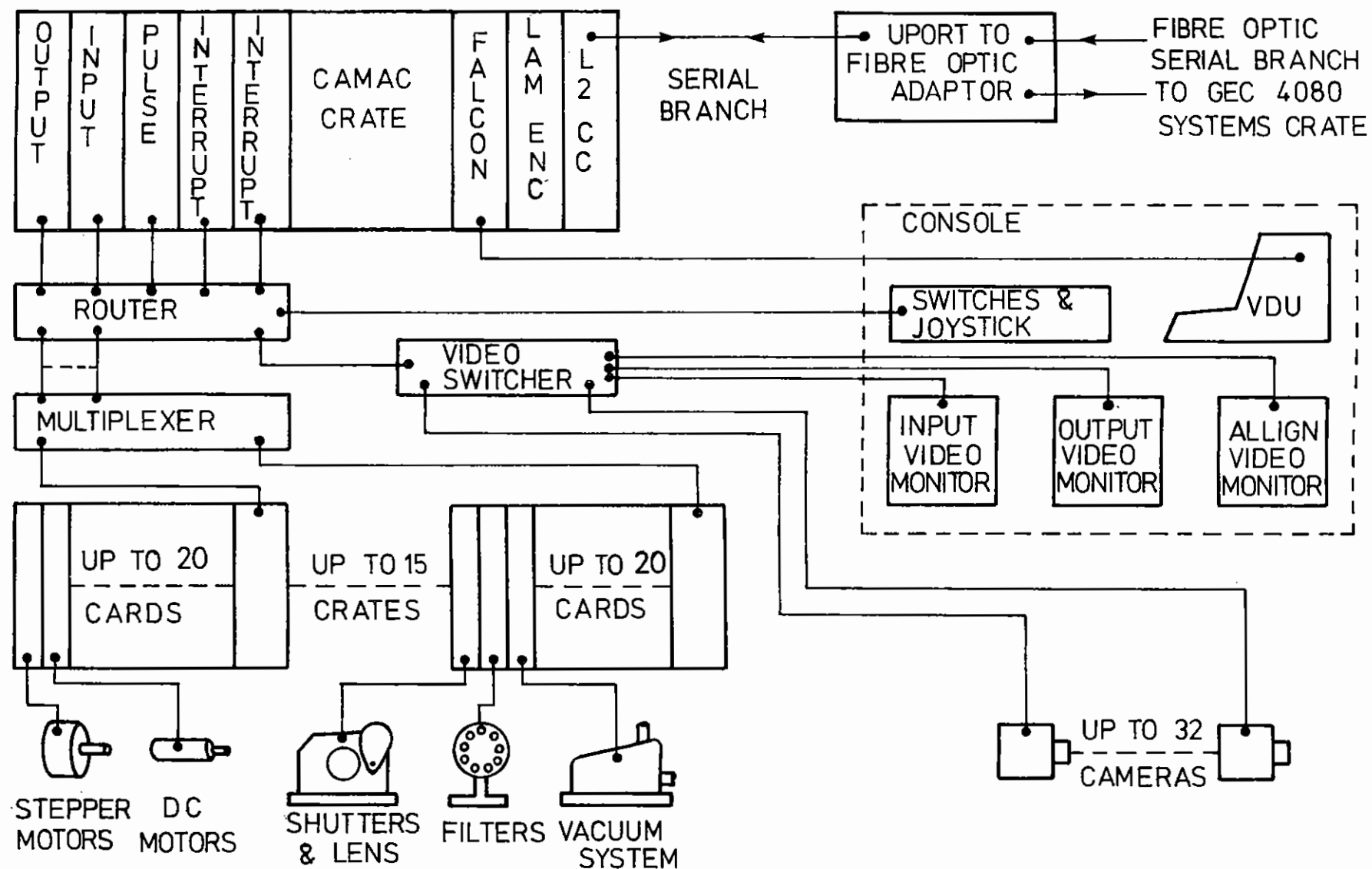


Fig. A6.21 THE TAW TARGET CHAMBER CONTROL SYSTEM

The buffer cards are connected to a Multiplexer which routes the information from the computer to the correct rack. It allows up to three racks to be accessed so that the x,y and z drives of a lens can be moved at the same time. All signals to the multiplexer originate from three single width CAMAC units—an output driver an input register and a pulse output unit. All actions performed by the device control system are entirely under software control.

A6.5.6 Software and Development System

C J Hooker, P Gottfeldt and C J Reason (RAL)

It was decided to write the application using a high level language in a real-time environment that would offer full concurrency thus allowing a modular approach to programming. DEC Micropower Pascal, offers these features. Programmes are developed on a system consisting of a VT103 terminal, LSI 11/23 with 128K bytes memory, RX02 with floppy disks, a printer, a PROM burner with the RT11 operating system and Micropower Pascal. To allow for debugging, a 64K RAM and housing were also purchased to attach to the Falcon.

The development cycle is as follows:

- a) The program is written in Micropower Pascal, compiled and the application built (including the Micropower Pascal Kernel) on the development system.
- b) The application is down line loaded into the RAM connected to the Falcon and debugged using the PASDBG program running in the development system.
- c) The program is then re-edited and the sequence repeated until the application is bug free.
- d) The application and the Kernel is then optimised and the debugging removed. It is rebuilt on the development system and blown into PROM.
- e) The RAM is removed from the Falcon and the PROM and bootstrap installed. The application then runs on power up.

The CAMAC driver is now written (called FALCAM) and is used by the test software to do CAMAC cycles and handle interrupts. The final application is now in the debug stage.

A6.5.7 Operator Interface

A R Damerell, M Forster, C J Hooker, C Reason and W T Toner (RAL)

The operators job is subdivided into tasks such as positioning targets, aligning beams on target, requesting and preparing for a shot etc. These are presented as a menu on the VDU and are selected by pressing one of a small set of special function keys which is physically separated from the rest of the keyboard. Secondary menus are then presented offering a choice of which beam to align etc. together with relevant information and instructions. When a choice is made at this level, control of the chosen set of devices is given to the joystick and switches on the console, slide mirrors and shutters are automatically set and the appropriate television station signals and controls are sent to the input (retro reflection), output (transmission) and align (reference comparator) TV monitors on the console.

Each monitor has a switch to select the near or far-field view and another to control brightness via a filter wheel at the camera station.

Limit switch stations and drive count are displayed on the VDU for the device which the operator has chosen to adjust via the joystick. When the task is complete the operator is prompted to declare whether (for example) a lens is aligned or not, and the programme retains this information for use in error checking. Warnings will be issued if the operator selects a beam which has already been aligned or if a shot is requested in a beam which has not been declared to be aligned, for example. In this way the operator is relieved as much as is possible of the burden of remembering a complex sequence of actions and is free to concentrate on matters involving judgement and precision alignment.

He is nevertheless free to ignore warnings or to select and control any

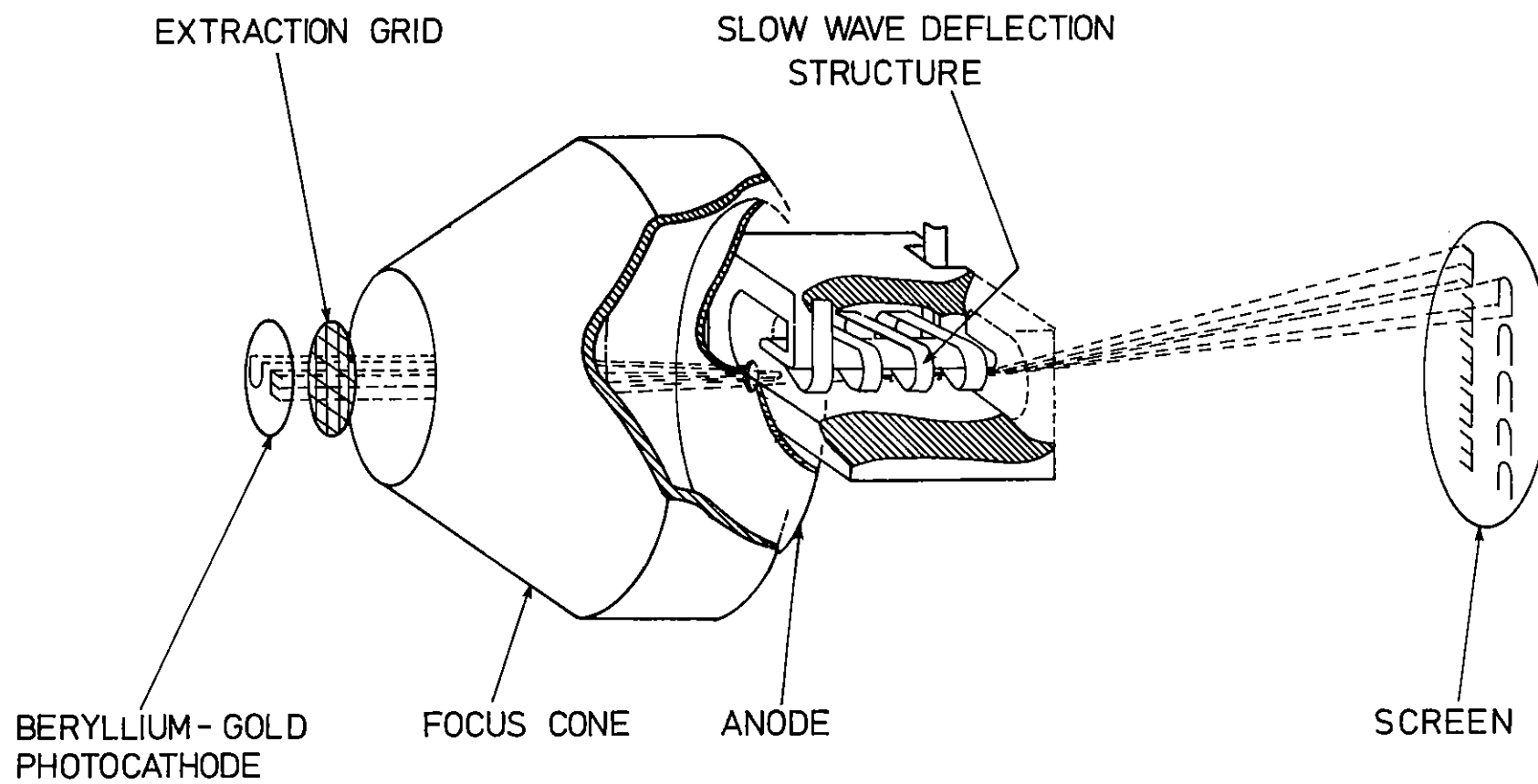


Fig. A6.22 Schematic of X-ray framing camera.

individual device in any beam lines out of sequence.

The Mark I version of the programme is limited to the control of the vacuum and alignment systems and permanent equipment protection shutters. It is intended to expand it later to include experimental devices. New equipment purchases will be computer compatible wherever possible. For example, a modular high voltage power supply was purchased last year for which a CAMAC controller will be bought so that device voltages may be programmed on before shots depending on the status of the vacuum system and off thereafter. It will also be possible to transfer selected status information and the operators response to prompts about target data and alignment method to the GEC 4080 computer for inclusion in the shot record.

A6.6 Aspheric Lens Manufacture

J Boon, D Nicholas and I Ross (RAL)

Of the two f2.5 lens designs for the new l2 beam target system described in last years Annual Report, the doublet design was judged to be the more predictable system based on previous experience gained in the manufacture of the f1 and f1.5 doublet lens systems. This view was supported by the manufacturers based on the flatness and the amount of departure from the best fit sphere, which relate to difficulties in polishing and final figuring of the surface.

The manufacturing techniques for these aspheric lenses has continuously been modified and updated on the basis of experience gained, whilst in broad outline the steps of manufacture remain the same, that is, an initial diamond tool grinding (A6.1) of the aspheric profile from the best fit sphere, followed by hand polishing and figuring. Steps have been taken to improve both these manufacturing stages in a number of ways in order provide an acceptable intensity distribution at the target plane.

In the original profile grinding of the f1 and f1.5 lens there was evidence that some irregularities had been introduced. This has been overcome by controlling both the first and second order time differentials on the

grinding head. The residuals have now been reduced to about the ultimate in grinding machine tolerance.

A remaining gross imperfection is a control feature in the aspheric generation due to machine dwell. This cannot readily be removed due to the lack of positive feedback signal from the grinding machine at this stage. This particularly manifests itself at the centre of the lens due to the low relative rotational speed of the lens to the grinding wheel. In the latest lenses an attempt to overcome this fault has been made by leaving the central ~25mm diameter of the lens polished at the best fit sphere radius, which leaves ~5µ to be polished off the ground portion of the lens. This has improved the overall features, but presented difficulties at the interface between polished and ground surface. It is easier to remove material from a ground surface than a polished surface. This again is being corrected by matching the ground finish with a similar grey finish to the best fit sphere. Six satisfactory doublets have been produced by adopting these techniques.

A6.7 Diagnostic Instrumentation Development

A6.7.1 A Subnanosecond X-Ray Framing Camera

N Finn, T A Hall and E McGoldrick (Essex)

This last year has seen the continuing development of the X-ray framing camera (A6.2) culminating in its use in an experiment to study 1.05µm implosions (see section A2.4). The camera now works reliably and has been used with framing rates of up to three frames per nanosecond.

Figure A6.22 shows the schematic layout of the camera. The principal of the slow wave structure is that a voltage pulse, travelling at the speed of light, follows a path such that it always acts on the same part of the electron beam which, in our case, is travelling at a approximately one quarter of the speed of light. By applying a staircase voltage waveform to the structure the electron beam is deflected in a series of steps. This waveform is produced by illuminating a silicon photoconducting switch with a

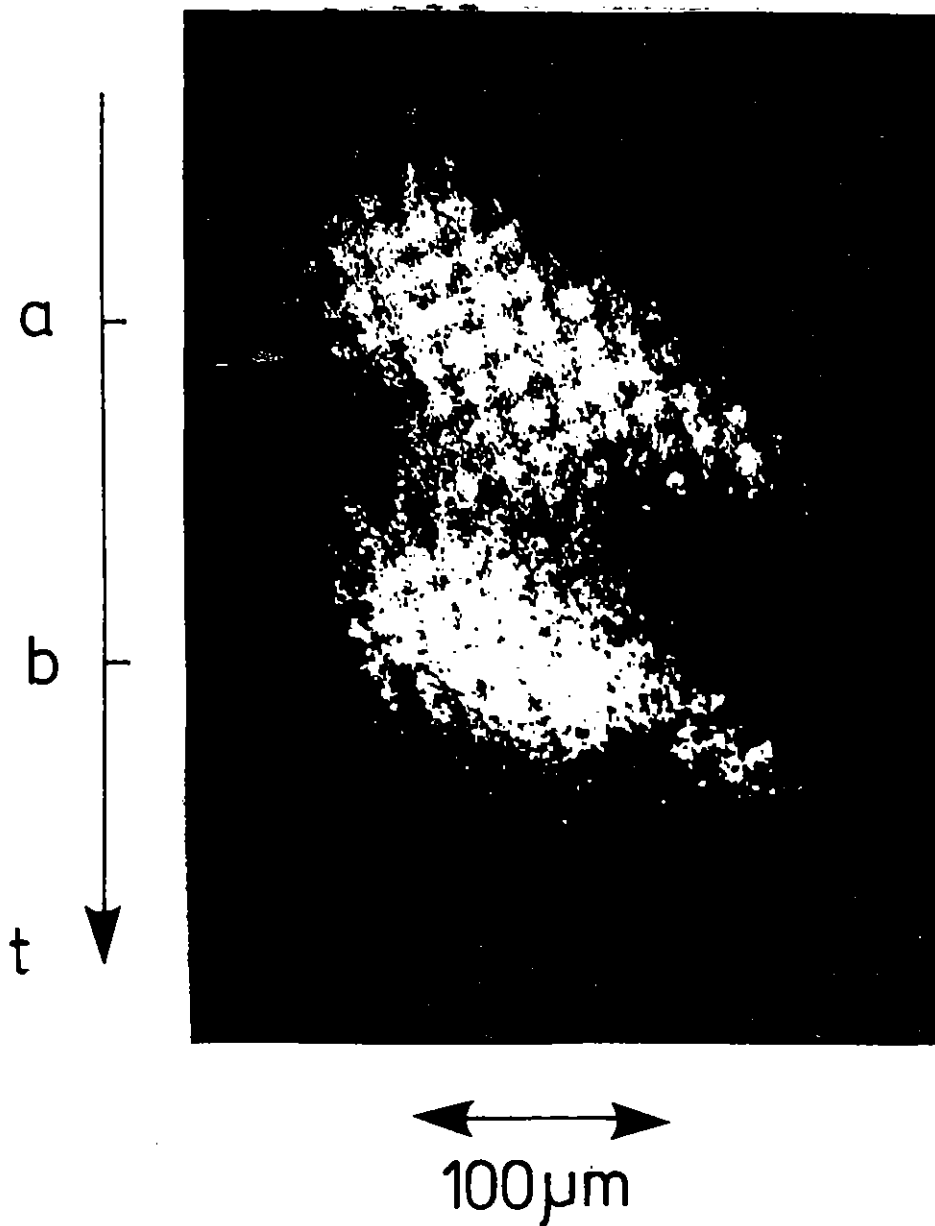


Fig. A6.23 Resulting images of resolution grid tests.

series of short, low intensity, pulses. Each pulse is of insufficient intensity to saturate the switch, so its conductivity increases in steps with each pulse.

The pulses are produced from a single pulse injected into a Fabry Perot etalon. Because successive pulses are of diminishing intensity the etalon is slightly mis-aligned, to separate the pulses in space, and then sent through a neutral density filter wedge arranged to give equal intensity pulses out. These are then focussed onto a diffusing screen in front of the switch.

Static image resolution tests have shown a cathode resolution of about 18 line pairs/mm. In dynamic tests at RAL a 2ns, 50J pulse was focussed onto a copper target backlighting a grid of 25μm periodicity with 10μm line and 15μm spacing. This was imaged onto the Caesium Iodide photocathode with X10 magnification, and framed at a rate of two frames per nanosecond. A 100ps, 70mJ pulse was passed through the etalon system to drive the silicon switch. The resultant images are shown in Figure A6.23. Only two images are seen because the backlighting x-rays were only sufficiently intense to be seen near the peak of the laser pulse. The first image is semi-static, the electron beam only moves from this position, while in the second image the beam moves into position and out. The difference in image quality can be seen, but the grid can still be resolved.

A new camera is currently under construction. The main changes will be an introduction of a blanking mechanism to "switch off" the electron beam between images - this will improve the quality of the image and lead to faster framing rates - and improvements will be made to the slow wave structure.

A6.7.2 Fully Automated Analysis of Alpha-particles and Proton Etched Tracks in CR-39 Plastic

A P Few (Bristol)

An image analysis system has been developed which is capable of fully

The system described here is completely software based, i.e. the hardware does no pre-processing of the image. The image is produced in a conventional light microscope and transmitted by a T.V. camera into a microcomputer. The image data is then transmitted in compressed form along a serial link to a VAX 11/750 computer where the main analysis takes place. The system operates with a microscope stage capable of stepping in 1 μm steps over 10 x 10 cm. The microscope is fitted with an autofocus system capable of positioning the objective to within 1 μm of the plastic surface. Typically, the system scans with an optical magnification of between 2.0 and 0.5 $\mu\text{m}/\text{pixel}$. At a magnification of 0.5 $\mu\text{m}/\text{pixel}$ the system can achieve an optimum resolution of 0.2 μm by interpolation along a line of pixels.

It is intended in the near future to acquire hardware to digitise the image frames directly into the analysis computer, thus providing a substantial increase in operating speed.

The image recognition and measurement of etched α -particle and proton tracks in CR-39 plastic is a particularly difficult problem because of the complex variation of track geometry with particle range, dip angle and etch conditions. In the ideal case there are five measurable parameters of the etched track: the major and minor axes of the opening mouth M_j, M_i , the total projected track length X_T , the track depth Z_T , and the diameter of the etched out end of the track, m . In the case of a non etched out track, $m = 0$. These parameters are indicated in Figure A6.24. In a given case, however, only a restricted number of these parameters will be accessible for measurement either in depth or in projection on the horizontal plane. Figure A6.25 shows some examples of how the etched track geometry may vary

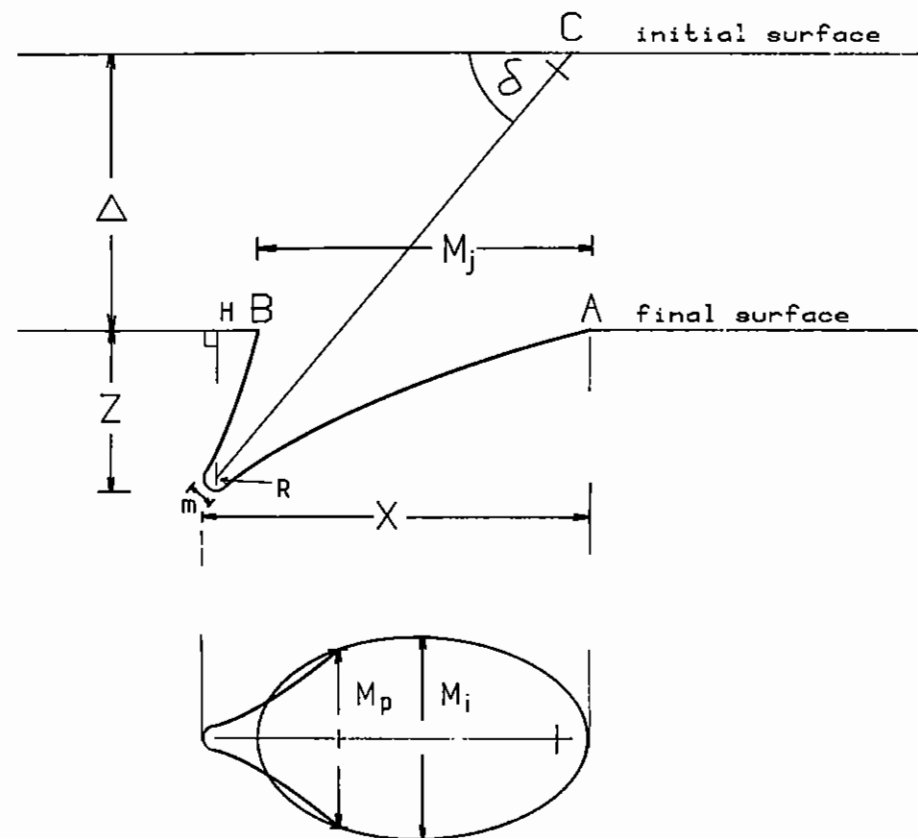


Fig. A6.24 A computer simulation illustrating the five measurable parameters of an etched α -particle track in CR-39. Note that Mp, the diameter of the mouth where the track trajectory intersects the post-etched surface is not a measurable parameter.

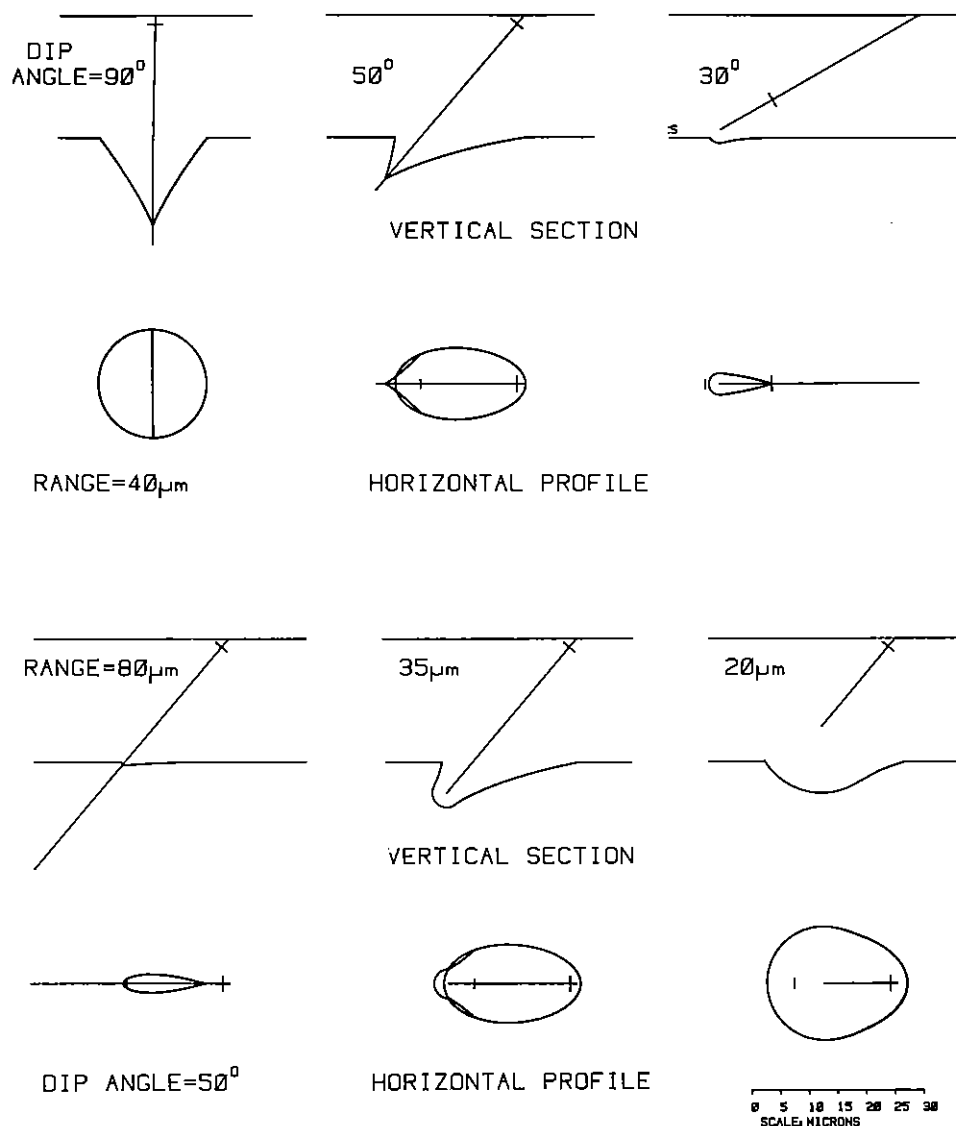


Fig. A6.25 Computer simulations illustrating the variation in etch-track structure with α -particle track geometry. They illustrate that in general only a restricted number of the parameters shown in Fig. A6.24 are available for measurement in the horizontal plane.

with parameters of the incoming particle and the etch conditions. The current system restricts the measurements to parameters projected on to the horizontal plane.

A variety of selection criteria have been developed which are applied in sequence to select efficiently tracks with well defined geometry and to reject non track defects on or in the plastic. In order to achieve the fastest frame analysis time, the selection criteria are applied throughout the measurement process: in general each calculational loop also contains a selection criteria integral. The selection criteria are:

- Minimum and maximum limits for X and Y size. This enables a crude first round selection to eliminate tracks that are considerably out of limits. To minimise computation time the limits are applied to X and Y directions only, and not to the longest axis.
- Area. The area of the event is computed and applied to selection limits.
- Mean optical density. The average grey level of the event is computed and compared to limits.
- Symmetry. The symmetry integral is evaluated over the area of the event and is applied to a cutoff limit.
- Convexity. The convexity of the event perimeter is computed and applied to a limit, so that a specified fraction of the perimeter that is concave will cause rejection of the event.
- Minimum and maximum limits for the Minor axis are next applied.
- Minimum and maximum limits for the X_i/N_j axis are next applied.

As indicated above a variety of possible track geometries exist with the result that in general only a restricted number of the parameters in the horizontal plane are available for measurement. The main categories are:

- M_i, X_T, m which applies to shallow tracks, $M_j < X_T$
- M_i, M_j which applies to steep tracks, $M_j > X_T$ or over etched tracks.

After the shape and position of the candidate track have been determined the program decides which parameter set to measure, either M_i, X_T, m or M_i, M_j .

It is envisaged that in future the Z drive on the microscope will be capable of being offset by known amounts from the surface focus position thus enabling depth measurements to be made, by maximising the sharpness at the deepest point of the track. However, it should be noted that the Z_T measurement itself will usually be of no value, since it will not have an accuracy comparable to that which is achievable on parameters measured in the horizontal plane.

The track analysis techniques will work with any arbitrary set of parameters. Therefore, any single parameter measured with an accuracy significantly less than the remainder will be rejected as unmeasured when the data set is passed from the image analysis program to the track analysis program.

Future developments will therefore be (i) to provide rapid multi-grey level input to the software (ii) calibrated Z_T drive from the surface focus position and (iii) development of the software to make the parameter measurements (both hidden and external) at calculated Z_T offsets from the surface focus.

The data from the image analysis software are passed to the track analysis program for computation of the input particle parameters. Depending upon the number of parameters given to the program it is possible to compute one or more of the following:

- 1) The particle range at the plastic surface, R
- 2) The particle dip angle with respect to the plastic surface,
- 3) The characteristic bulk etching rate of the plastic V_B
- 4) The track etch rate V_T . This may be reduced to the value of the track etch rate at a particle range of $10\mu\text{m}$, V_{T10} . This serves as a single parameter to scale the general form of the track etch rate versus particle range relationship
- 5) The etch induction time.

In practice only a restricted number of parameters are passed to the track

analysis program so that all five of the above quantities are rarely determined simultaneously.

Initial testing of the system has been carried out with ^{252}Cf α -particles and also with protons of energies up to 2 MeV from a Van de Graaf accelerator.

To illustrate the results obtained on α -particles, Figure A6.26 shows the measured distribution of end radii, $m/2$, which in practice is the most difficult parameter to measure. All the data shown were obtained for tracks where the system measured the parameters X_T , M_1 and m . The measurements were made at a magnification of $0.5\mu\text{m}/\text{pixel}$ and the instrument resolution is estimated to be $\sim 0.3\mu\text{m}$. The spread in the main signal, between $\sim 3.4 - 4.4\mu\text{m}$ is chiefly due to real differences in the range of the α -tracks. Measurements less than $3\mu\text{m}$ arise principally from etch pits blocked by etch products.

The data from the image analysis system were then analysed by the method described above to fit a bulk etch rate, V_B , the track etch rate scale factor V_{T10} and a total range to each track individually. The resultant V_B distribution is shown in Figure A6.27, the V_{T10} distribution is shown in Figure A6.28 and the range distribution in Figure A6.29. All three graphs contain all the measured data, and no selection was performed to improve the distributions. It can be seen that on individual tracks the bulk etch rate can be determined to better than 10%, and V_{T10} to $\sim 4\%$. The mean range of the tracks in Figure A6.29 is $35.3\mu\text{m}$ with a standard deviation of $0.6\mu\text{m}$, which is marginally better than a skilled human observer could achieve using the same microscope objective and measuring only X_T , M_1 and m . Improvements to the system hardware, as described above, will be required before it will be possible to use a higher resolution objective.

Figure A6.30 illustrates the measured distribution of minor axes of normally incident 2 MeV proton tracks. The spread in distribution (σ) is $\sim 0.22\mu\text{m}$ which is close to the theoretical optimum value achievable with a magnification $0.5\mu\text{m}/\text{pixel}$. This would correspond to an uncertainty in the

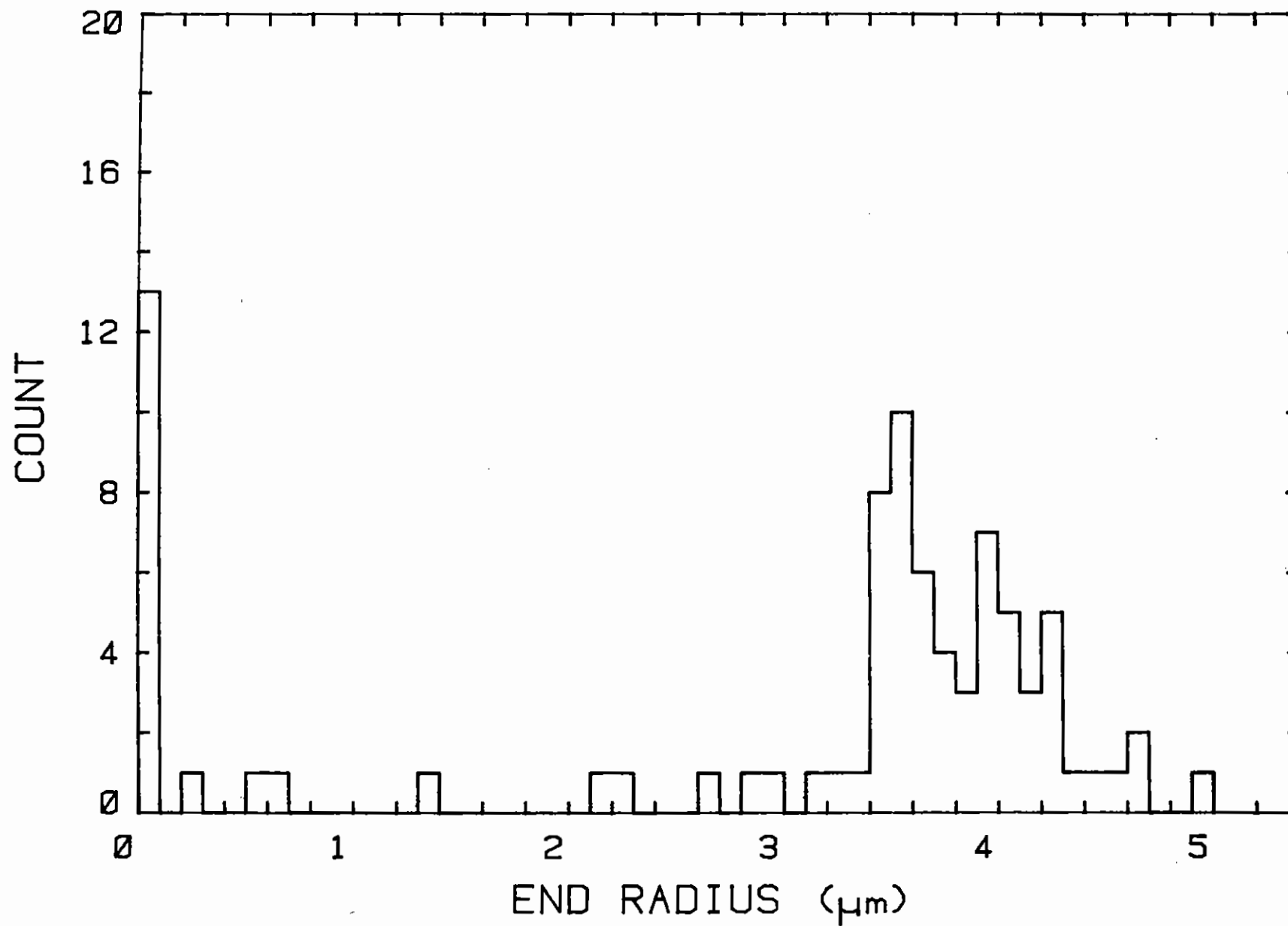


Fig. A6.26 Distribution of the measured values of end radius, $m/2$ on etched out ^{252}Cf α -particle tracks in CR-39.

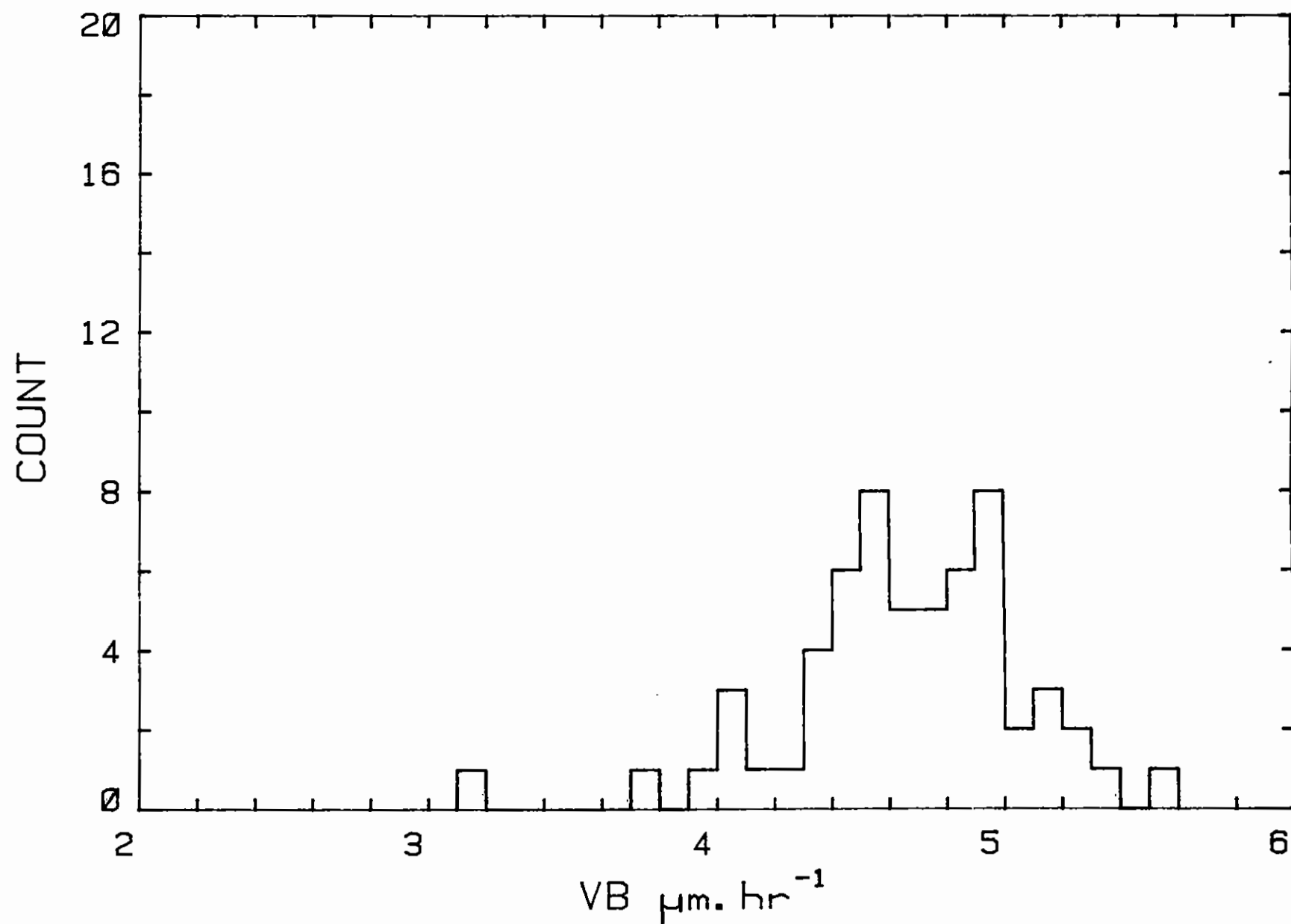


Fig. A6.27 Distribution of measured values of bulk etch rate, V_B from automated scanning measurements of ^{252}Cf α -particle tracks in CR-39.

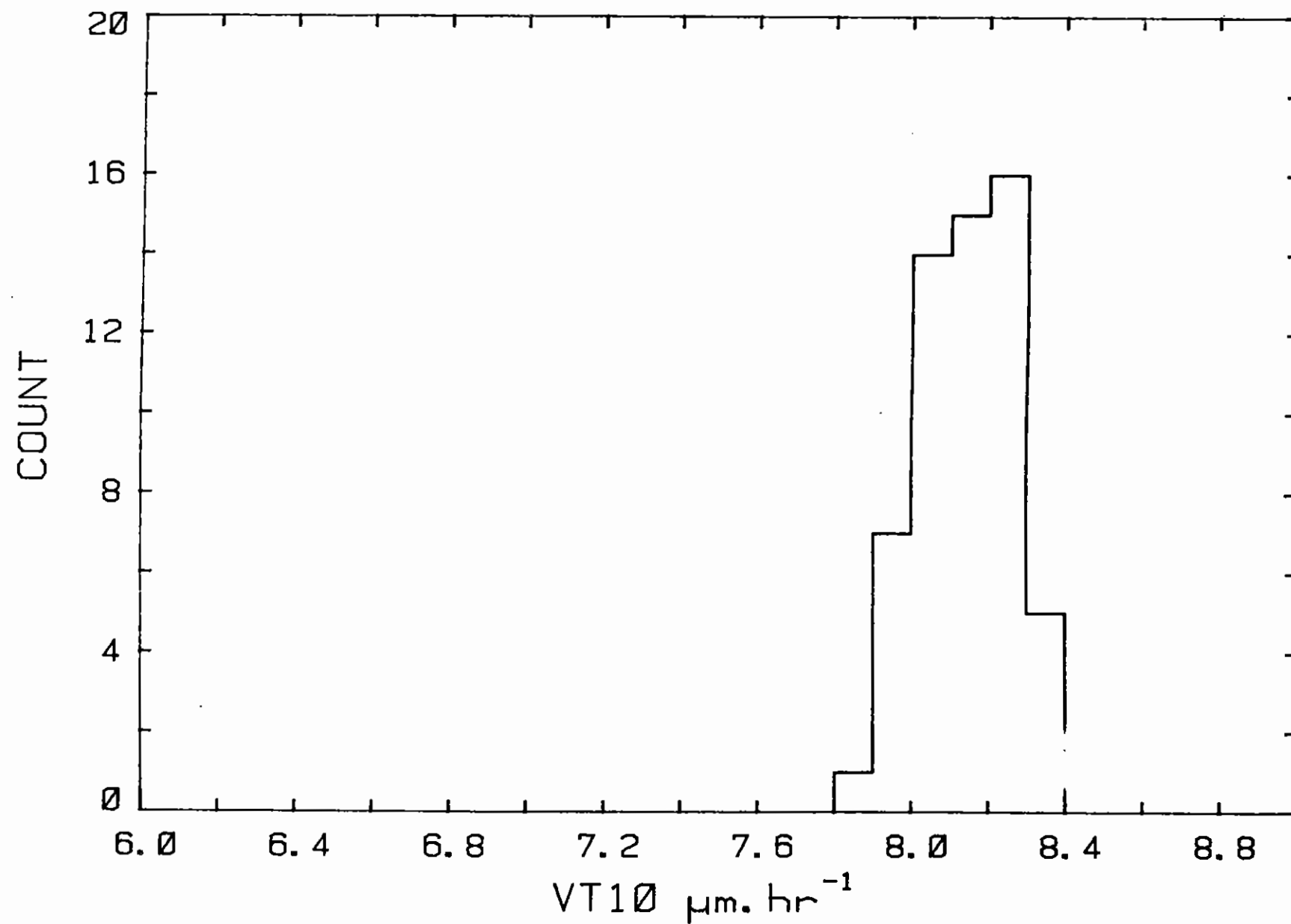


Fig. A6.28 Distribution of measured values of the track etch rate scale factor V_{T10} .

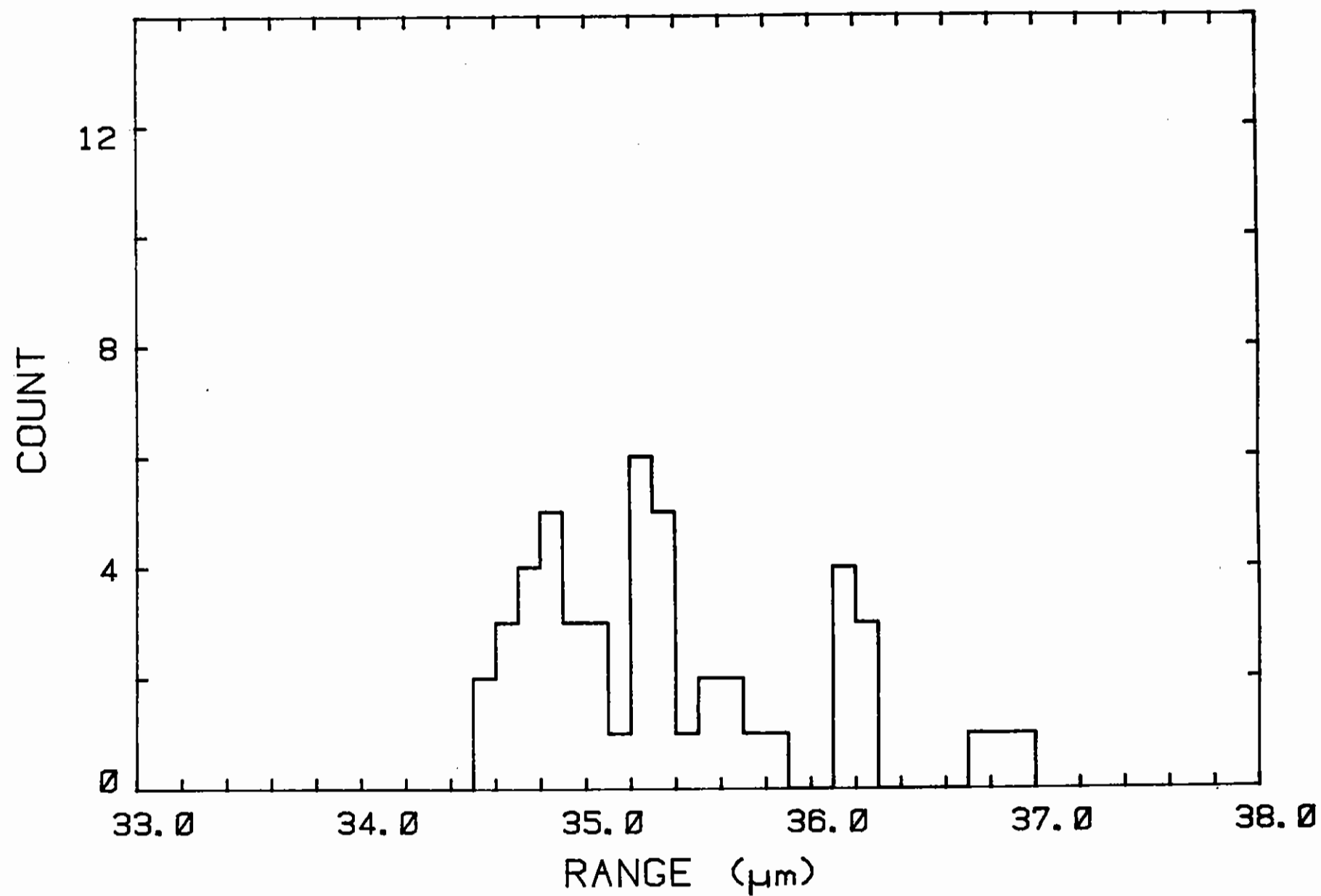


Fig. A6.29 The range distribution of ^{252}Cf α -particle tracks in CR-39 as determined by automated image scanning and analysis.

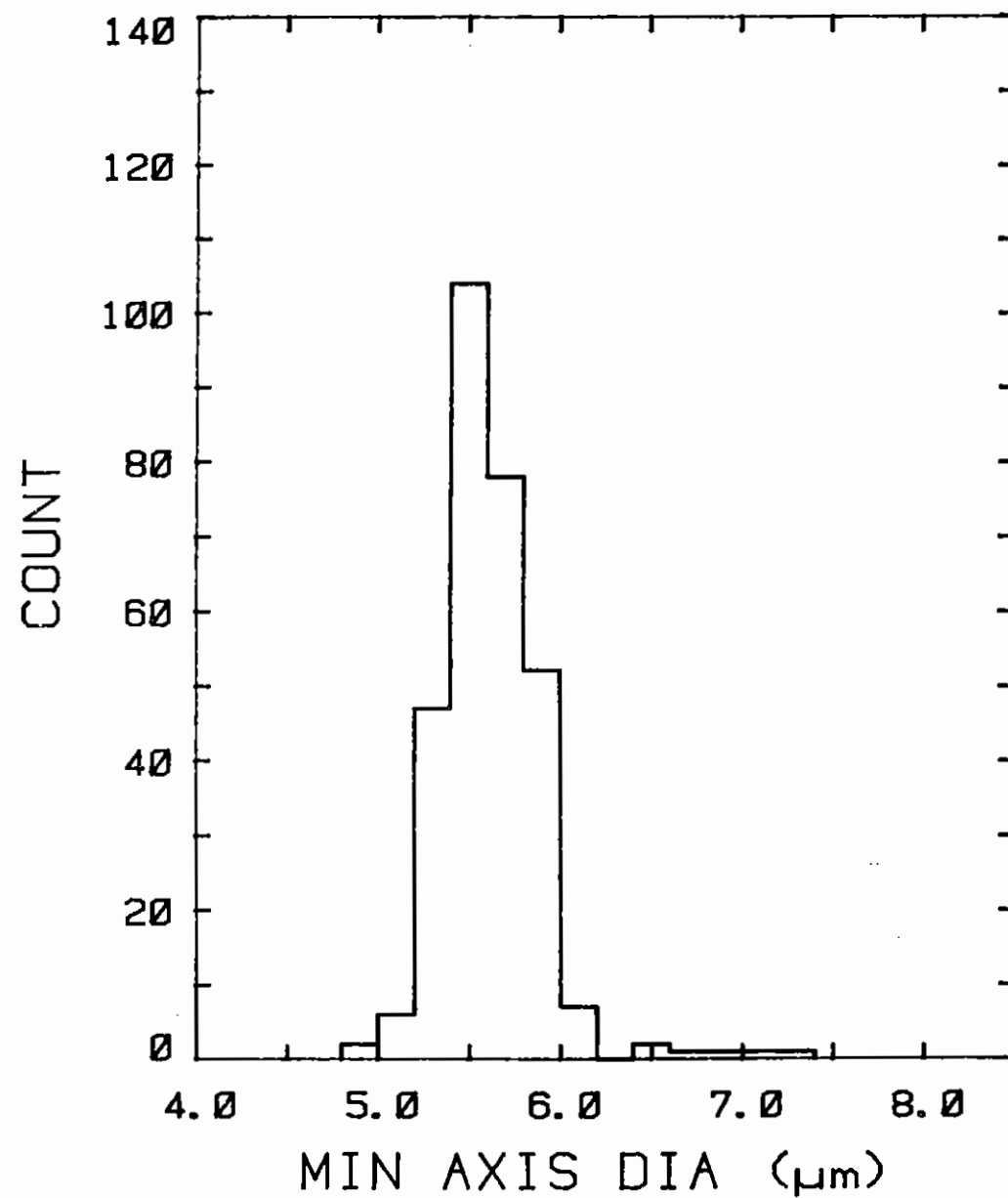


Fig. A6.30 The distribution of minor axis diameter M_1 on 2 MeV proton tracks at normal incidence in CR-39.

proton range of 5 μm , corresponding to an energy uncertainty of 0.1 MeV. The CR-39 plastic was in this case etched for a bulk removal, $\Delta \sim 8 \mu\text{m}$. The energy resolution could be improved by a factor ~ 4 -5, by etching the plastic for the optimum amount of surface removal, $\Delta \sim 30 \mu\text{m}$ in this case.

Apart from the analysis of thermonuclear particles the analysis system has wide application for nuclear particle analysis in other fields. In particular the system is being developed for application in the growing areas of research into the uptake and retention of α -radionuclides in man and in neutron spectrometry and dosimetry.

A6.7.3 Compact X-ray Streak Camera

J Hares (IC)

The compact x-ray streak camera described in the 1983 annual report (Section 1.7.3) has been used on several experiments in this period. These include colliding shock absorption spectroscopy, planar dot spectroscopy and time resolved EXAFS. In particular the planar dot experiment has benefited from the small physical size in that a doubly reentrant part arrangement (see Figure A6.31) was used to get the photocathode to a distance of only 9cm from the plasma. This resulted in a very high sensitivity.

The recombination continuum from a 50nm thick dot, 70 μm in diameter, was recorded out into the coronal region. Hydrogen and helium like line ratios were obtained from 25nm thick dots. The dot, being very thin, samples a small region of the ablation plasma as it travels outwards, thus time translate into Z a the streak record.

The solid state sweep circuitry has been totally reliable.

A6.7.4 A Gated X-Ray Intensifier with a Resolution of 50 Picoseconds

A K L Dymoke-Bradshaw, J D Kilkenny, J Westlake (IC)

Many aspects of laser produced plasmas and laser fusion studies require

time-resolved X-ray images, in particular the study of implosion symmetry and Rayleigh Taylor instabilities. The useful exposure times required are determined by the velocity, 10^5ms^{-1} and the 10 μm resolution to be 100ps or less. A gated X-ray intensifier could ultimately produce exposure time of 50 ps.

The device, described in the annual report last year, was a proximity focussed tube with photocathode, extraction mesh, microchannel plate and phosphor. This device suffered from a poor contrast ratio as X-rays passing through the photocathode formed an image superimposed on the gated image.

To improve the contrast ratio we have built a second tube in which the gated and ungated images are separated by drifting the photoelectrons for 35mm at 12° to the straight through X-rays. The photoelectron image is drifted in a magnetic field of 0.2 to 0.3 T which gives an electron Larmor radius of $\sim 25\mu\text{m}$ in the direction perpendicular to the drift direction. The tube is shown in Figure A6.32. Electrostatic imaging of the photocathode onto the micro-channel plate was precluded as the large longitudinal energy spread of the gated electrons would result in severe chromatic aberration. The elements of the tube remain essentially unchanged from last year except for the increased separation of the mesh and micro-channel plate.

With a large magnetic field on the tube the criteria for proximity focussing in the photocathode-mesh region can be relaxed. It is no longer necessary that the electrons cross the whole photocathode to mesh distance during the gate pulse as long as they pick up sufficient energy to make them detectable by the micro channel plate ($\sim$ two hundred eV) and sufficient energy to overcome the reverse bias applied to the photocathode. Also resolution is fixed by the Larmor radius not the proximity focus point spread function. This means that the photocathode to mesh distance and the width of the stripline can be increased giving more field of view without departing from 50 ohm geometry. The present device uses a 1.5mm spacing with a photocathode 7mm wide by 10mm long.

Static resolution tests using a pulsed (15ms risetime) 0.3 Tesla field were performed. A continuous microfocus beam target X-ray source was used, run

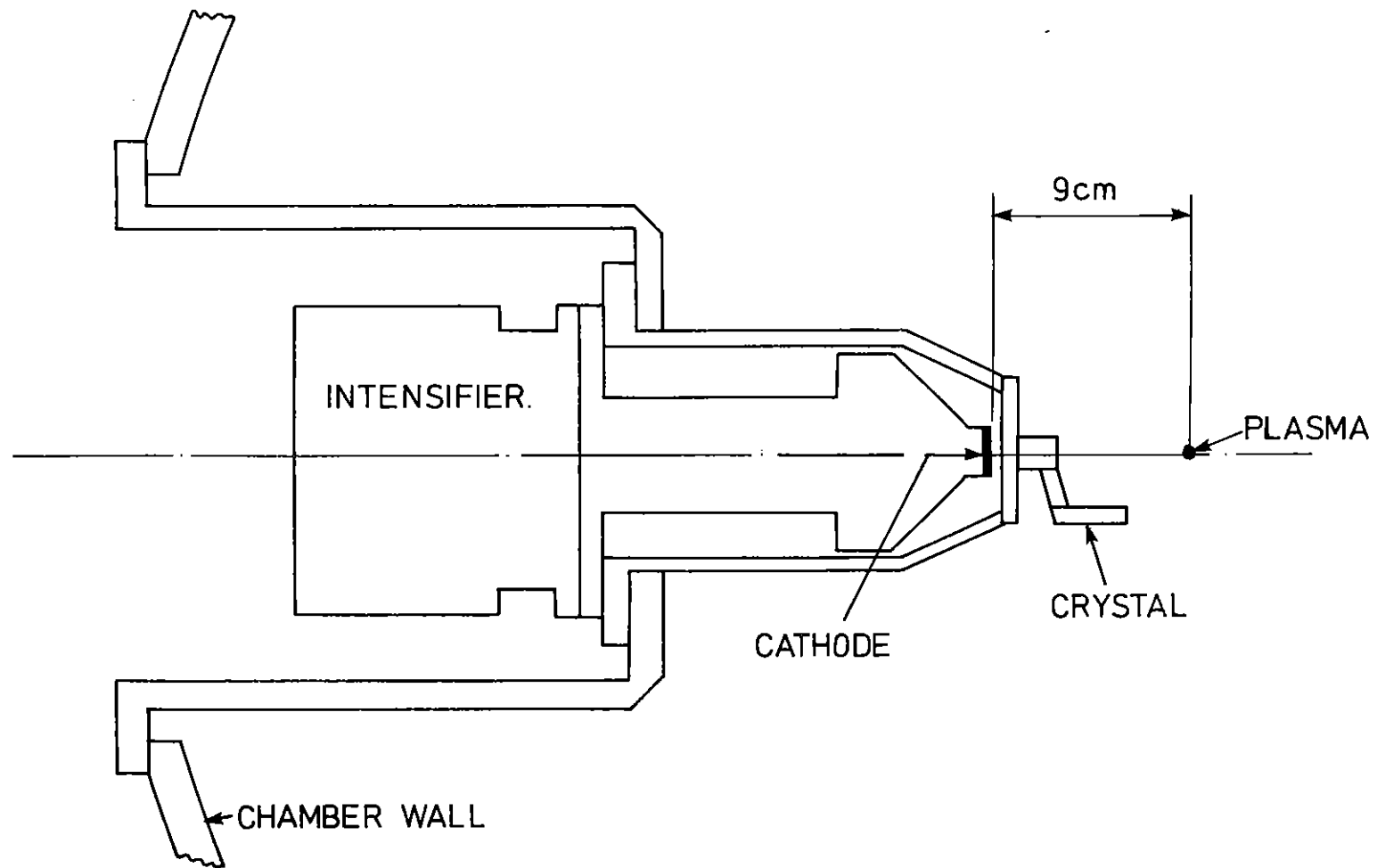
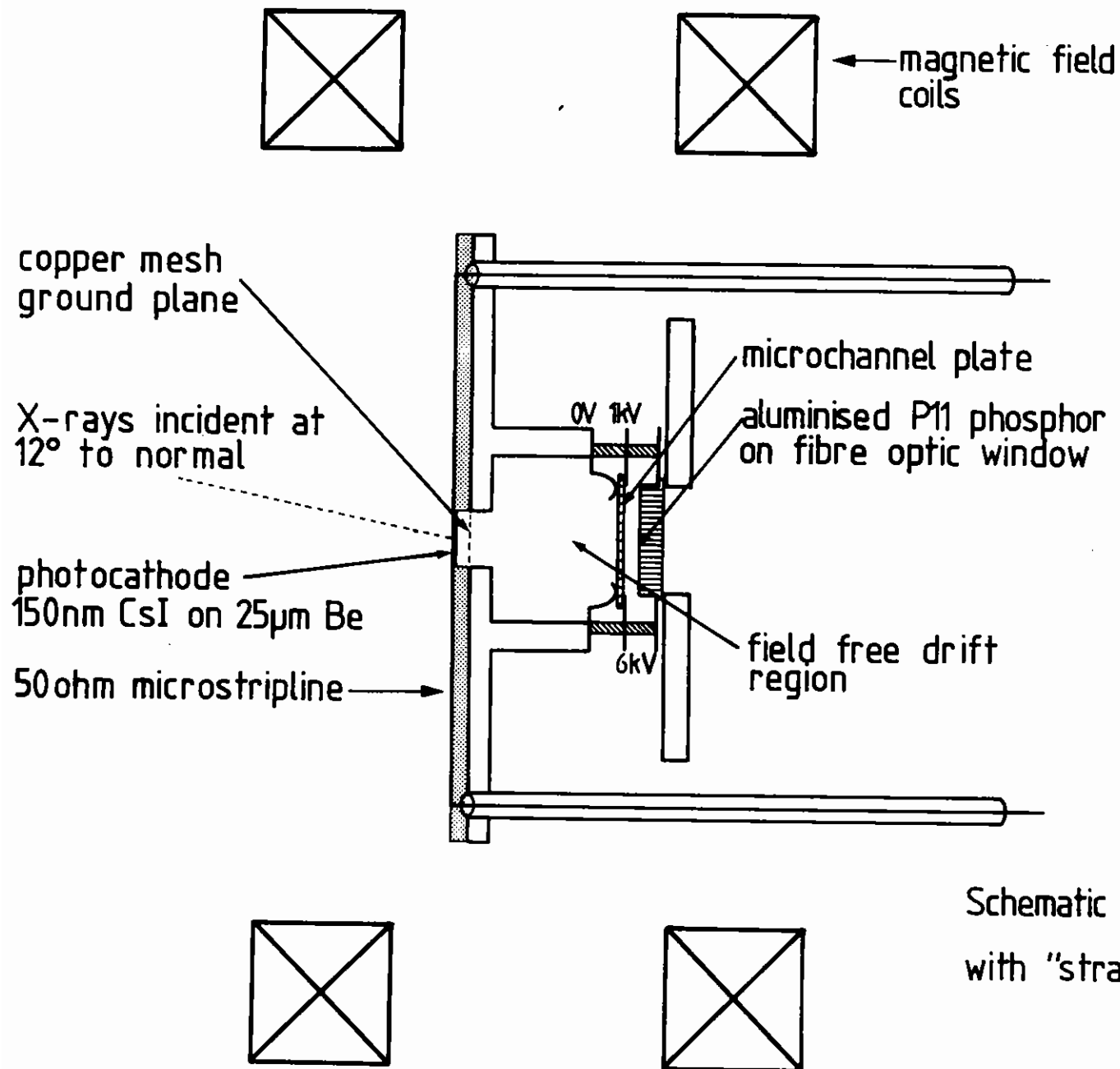


Fig. A6.31 Illustrating doubly re-entrant geometry for high sensitivity.



Schematic of new image tube
with "straight-through" image
rejection

with 6kV electrons onto an aluminium target. The photocathode was gated on for 500 μ s at the peak of the magnetic field. The H.T. for the microchannel plate and phosphor was gated on for about 50ms.

To measure the spatial resolution several copper meshes were placed on the front of the photocathode. The results shown in Figure A6.33 are from the superposition of 40 exposures. It can be seen that the 80 μ m pitch 1:1 mesh is easily resolved giving 12.5 line pairs per mm. The mesh with holes of size 230 by 23 μ m also shows some signs of being resolved in the smaller dimension. With softer X-rays (or thicker meshes) and a higher magnetic field it is expected that 20 line pairs per mm should be possible.

It can also be seen in Figure A6.33 that the bottom of the picture is a duplicate of the top. This is the straight through image which has been exposed for one hundred times longer.

Gating tests on the tube were performed on the Vulcan laser at the Rutherford Appleton Laboratory. A laser pulse of several joules and 120 ps duration produced an X-ray source on each side of a gold on Mylar foil target. The sources were separated by about 400 μ m in space and 100 picoseconds in time. The sources were imaged onto the photocathode with a 100 μ m pinhole which gave sufficient resolution to separate the sources. Part of the laser beam was used to drive the Auston switch. The gate pulse length was set at 200 ps and on a shot to shot basis the gate pulse timing was scanned over a nanosecond.

The results are shown in Figure A6.34. By gating very early neither source is imaged. As the gate timing is made later, first one then both of the X-ray sources is imaged. The sources have fast risetimes but slow decay times.

It can be seen that a change of 50 picoseconds in the timing can change the resulting image from barely detectable to well exposed.

The contrast ratio appears to be very high and we have yet to measure it. Emission of photoelectrons from the mesh which could reduce the contrast

ratio does not appear to be a severe problem.

Having produced a device of useful temporal and spatial resolution with good gating contrast we next intend to build a multichannel device with several photocathodes and one intensifier. Imaging will be performed with pinhole arrays or a Kirkpatrick Baez microscope. An Auston switch pulser will be required for each channel. A larger magnetic field of up to 1T may be used to improve spatial resolution so that it is MCP limited.

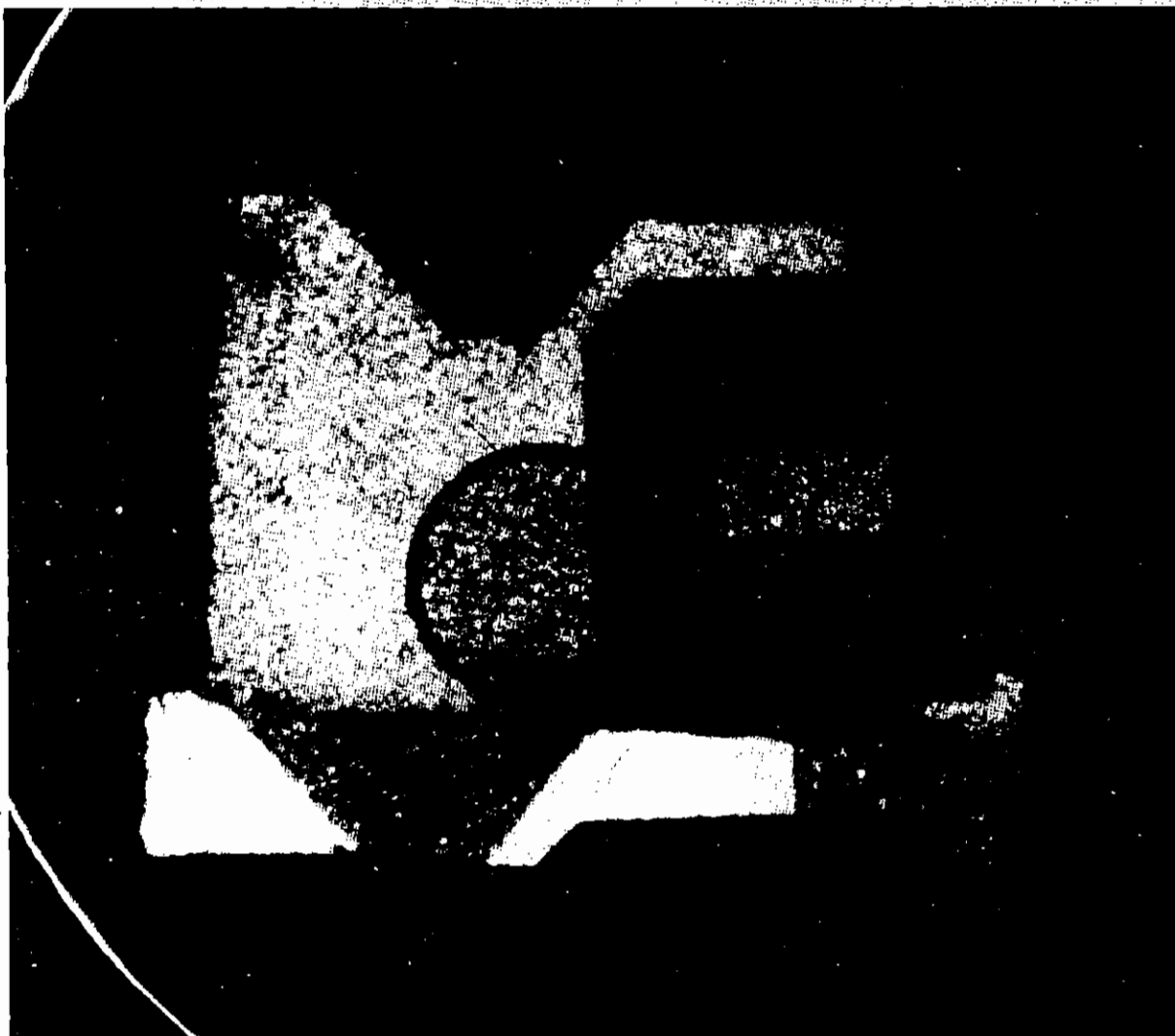
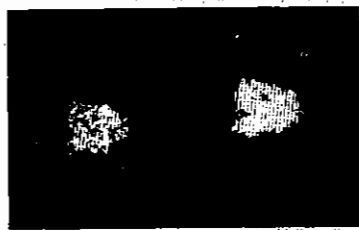


Fig. A6.33

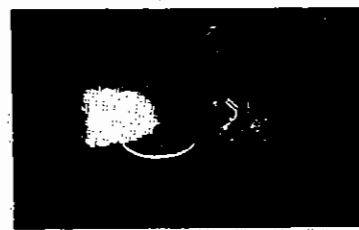
Static resolution tests. The superposition of 40 half millisecond exposures. The photocathode was masked with two grids. On the right is 1:1 grid with 80 micron pitch showing resolution at 12.5 line pairs per millimetre. To the left is a mesh with rectangular holes of size 230 by 23 microns. At the bottom of the picture can be seen a duplicate of the image at the top. This is the straight through image and is the superposition of 40 fifty millisecond exposures.



150 ps



300 ps



750 ps



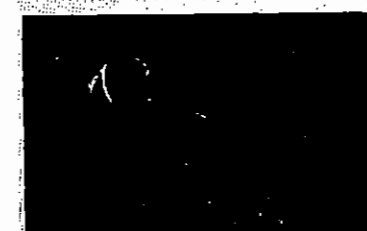
200 ps



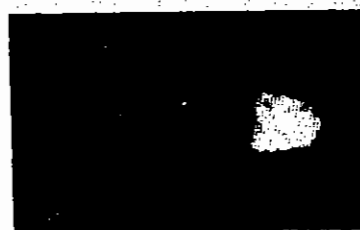
500 ps



850 ps



1050 ps



250 ps



650 ps



950 ps

IMAGES OF 2 PLASMA X-RAY SOURCES TAKEN WITH THE GATED X-RAY INTENSIFIER

Fig. A6.34

Time resolution tests with the second tube. Two X-ray sources were produced either side of a gold on mylar foil and imaged (with low spatial resolution) onto the photocathode. The laser beams had a relative delay of about 200 picoseconds. It can be seen that the left spot comes on later than the right spot, and turns on in about 50 picoseconds. It can be seen that the extinction ration is now excellent.

References

- A6.1 D J Nicholas and J E Boon, J Phys D: Appl Phys 14, 593 (1981)
- A6.2 Annual Report to the CLF, Rutherford Lab Report RL-83-043 (1983)

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Editor M C Gower

SECTION B ULTRA VIOLET RADIATION FACILITY SCIENTIFIC PROGRAMME

INTRODUCTION

M C Gower (RAL)

The past year was the second year of operation of the Ultraviolet Radiation Facility (UVRF) in the CLF, and was one of intense activity by the researchers who made use of the facility. During the year a total of 51 users drawn from 16 university groups used the facility. In addition, the facility was hired, by the purchase of time on the facility, by three other establishments.

The primary laser sources of the facility are a line-narrowed high brightness excimer system for tunable operation around wavelengths of 193, 249 and 351nm, an excimer pumped dye laser continuously tunable between 320 and 980nm and a high energy F_2 laser operating in the vacuum ultraviolet spectral region at 157nm. A range of various sophisticated diagnostic and detection equipment is also available in the facility.

Scheduling meetings to consider and allocate time on the facility were held in July and November 1983 and March 1984. In contrast to last year, three scheduling meetings are now held every year so that four (instead of six) monthly periods are scheduled at each meeting. As last year, bids for time on the facility greatly exceeded the time available. In response to this increasing demand, in September 1983 the facility expanded into an additional laboratory adjacent to the original UVRF lab. Experiments which only require use of the dye laser or the line-narrowed excimer laser system can now be scheduled independently for simultaneous operation. For those experiments which require the use of both of these systems it is possible to operate both laboratories in tandem and relay beams between them. As a result of this expansion the number of experiments which have been scheduled has significantly increased.

The bids for Division Head's weeks in both areas have also been greater than the time available. However, by use of more flexible periods at weekends,

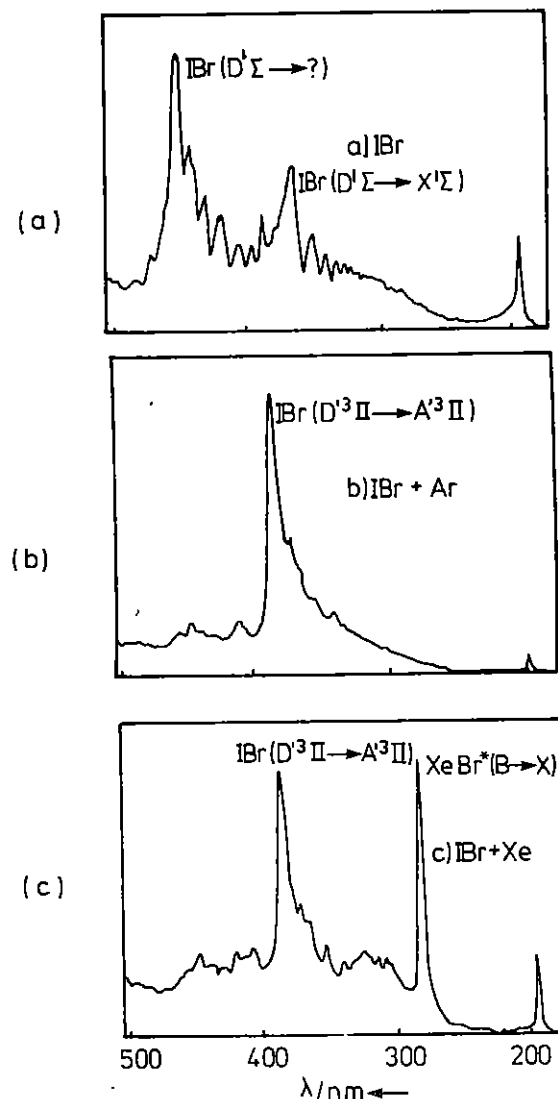


Fig. B1.1 Fluorescence observed following excitation of IBr at 193 nm (ArF laser)

(a) IBr ($P_{\text{IBr}} = 173 \text{ Nm}^{-2}$, $P_{\text{Br}_2} = 176 \text{ Nm}^{-2}$)

(b) IBr + Ar ($P_{\text{IBr}} = 160 \text{ Nm}^{-2}$, $P_{\text{Br}_2} = 147 \text{ Nm}^{-2}$, $P_{\text{Ar}} = 7.45 \text{ k Nm}^{-2}$)

(c) IBr + Xe ($P_{\text{IBr}} = 173 \text{ Nm}^{-2}$, $P_{\text{Br}_2} = 176 \text{ Nm}^{-2}$, $P_{\text{Xe}} = 7.09 \text{ k Nm}^{-2}$)

evenings and time set aside for maintenance it has been possible to accommodate most of those users who wished to conduct short exploratory experiments.

Throughout the year the equipment in the facility has been expanded and enhanced. In particular, a Lambda-Physik Model EMG150ETS has been installed to replace the EMG101E/201 oscillator/amplifier configuration used for line-narrowed operation. A high energy (40mJ) F_2 laser (Oxford Lasers KX-1) which operates at 157nm has also been installed. Additional diagnostic equipment has also been added to the facility.

As last year, the main activities in the facility were photochemistry (61% of scheduled time) and physics (28%) related activities. However, photobiology (4%) and engineering (2%) experiments were also carried out. It can be seen from the work described below that the year has been extremely productive for the users of the UVRF.

B1 PHOTOCHEMISTRY AND SPECTROSCOPY

B1.1 Photochemical and Spectroscopic Studies of Small Molecules

R J Donovan, M MacDonald, C McKendrick, A Torre, J P T Wilkinson (Edinburgh), C Fotakis (University of Crete) and M Gower (RAL)

During the past year a number of new projects have been undertaken together with further work on some of the projects described in last years report. Particular emphasis this year has been placed on extending our work on the spectroscopy and reactions of the halogen ion-pair states.

B1.1.1 Oscillatory Continuum Emission from IBr

Two strong oscillatory continuum systems were observed (B1.1) following excitation of IBr at 193 nm (ArF laser). As IBr exists in equilibrium with both I_2 and Br_2 and because I_2 gives intense fluorescence when excited at 193 nm, a small excess of Br_2 was added to all samples. By careful choice of conditions, good quality spectra were obtained (see Figure B1.1 below).

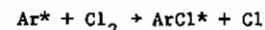
This work is described in further detail in reference B1.1. We are currently in the process of analysing these spectra to obtain new spectroscopic data on excited states of IBr.

B1.1.2 Oscillatory Continuum Emission from Br_2

Last year we reported that three new oscillatory continuum systems had been observed following excitation of Br_2 at 157 nm (F_2 laser) (B1.2). Further work has been carried out and the quality of the spectra has been greatly improved. We are currently analysing these spectra using the Edinburgh computer code.

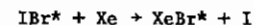
B1.1.3 Reaction of Excited IBr with Xenon

The reactions of electronically excited rare gas atoms with halogens and halogen containing molecules, e.g.

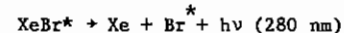


have been studied extensively due to the importance of excited rare gas halide molecules in the development of high power laser systems. It has also been shown that electronically excited Cl_2 reacts with krypton to yield electronically excited $KrCl^*$ (B1.3).

We have now shown (B1.4) that IBr excited at 193nm (ArF laser) reacts with ground state xenon atoms to yield excited $XeBr^*$, viz.:



Fluorescence from excited $XeBr^*$ is observed at 280 nm (see Figure B1.1):



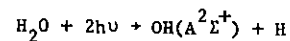
We have also studied the collisional quenching of excited IBr by Ar which populates the lowest ion-pair (D') state giving strong fluorescence at 385 nm. Further details are given in reference B1.4. From this and earlier work it appears that it will be possible to construct a range of new and efficient optically pumped lasers involving halogen and interhalogen molecules.

B1.1.4 Two Photon UV Excitation of SO₂ and Laser Induced Fluorescence from SO

We have shown that strong fluorescence from $\text{SO}(\text{B}^3\Sigma^- \text{ } (v' = 2) + \text{X}^3\Sigma^- \text{ } (v'' = n))$ is observed following KrF laser (248nm) excitation of SO_2 . The SO radicals are initially formed in their ground electronic state by a sequential two-photon absorption mechanism involving the $\tilde{\text{B}}^1\text{B}_1$ state of SO_2 as an intermediate state and a high-lying Rydberg state as the dissociative state. The SO radicals then absorb a further photon at 248 nm which gives rise to the observed $\text{SO}(\text{B-X})$ fluorescence (see Figure B1.2). This work is described in further detail in reference B1.5.

B1.1.5 Two Photon UV Excitation of H₂O

Work on the two photon excitation of H_2O at 248 nm using the line narrowed output from a KrF laser has recently been carried out. Fluorescence from $\text{OH}(\text{A}^2\Sigma^+)$ has been monitored as the output from the line narrowed system was tuned over the gain profile and considerable structure in the excitation spectrum has been observed.



The spectroscopic data from these experiments is currently being analysed.

B1.2 Photochemistry and Raman Spectroscopy

R E Hester, P M Killough, J E Bennett, W E L Grossman (York)
R Wilbrandt (Risø National Laboratory, Denmark)

B1.2.1 Introduction

Our work over the past year has been of two types:

(i) time-resolved resonance Raman (TR^3) spectroscopy of photochemical reaction intermediates, and (ii) resonance Raman spectroscopic studies with ultraviolet excitation. Both types of experiments depend on the versatile

excimer laser system and on the sensitivity of the diode array multichannel detection system available at the RAL UVRF.

TR^3 experiments exploit the pulsed nature of the excimer and dye lasers. An initial photolysis or "pump" initiates the photochemistry and a subsequent probe pulse interrogates the system. The probe pulse may be separated in time and/or in wavelength from the pump using our present experimental arrangement (as previously reported). Thus it is possible to monitor the progress of a photochemical reaction as a function of time and at a convenient (usually a resonance) wavelength. In principle TR^3 offers more information than pulsed absorption spectroscopy (flash photolysis) since the results are the highly structured vibrational spectra of the species in the solution, rather than broad featureless absorption spectra.

The ability to excite Raman and resonance Raman spectra in the ultraviolet is important to the practice of Raman spectroscopy because the intensity of the scattered light is proportional to the fourth power of the frequency of the exciting laser light. It is especially important for resonance Raman since the majority of substances have their most intense absorption bands in the ultraviolet region. Relatively few compounds are coloured and thus accessible to resonance enhancement using visible lasers. Since the resonance enhancement effect is specific to the chromophoric part of the molecule, it may be used to good advantage for selective probing of small structural groups in large molecules (eg. enzyme active centres) with good sensitivity.

B1.2.2 Experimental Methods

Since our last report, several modifications have been made to our experimental arrangement. The most important of these has been the addition of a PARC model 1211 high voltage pulser to gate on and off the intensifier in the multichannel detector. This significantly improves the quality of the spectra obtained because the initial photolysis pulse and any associated fluorescence are gated out.

Also of importance to our work has been the greater reliability and improved

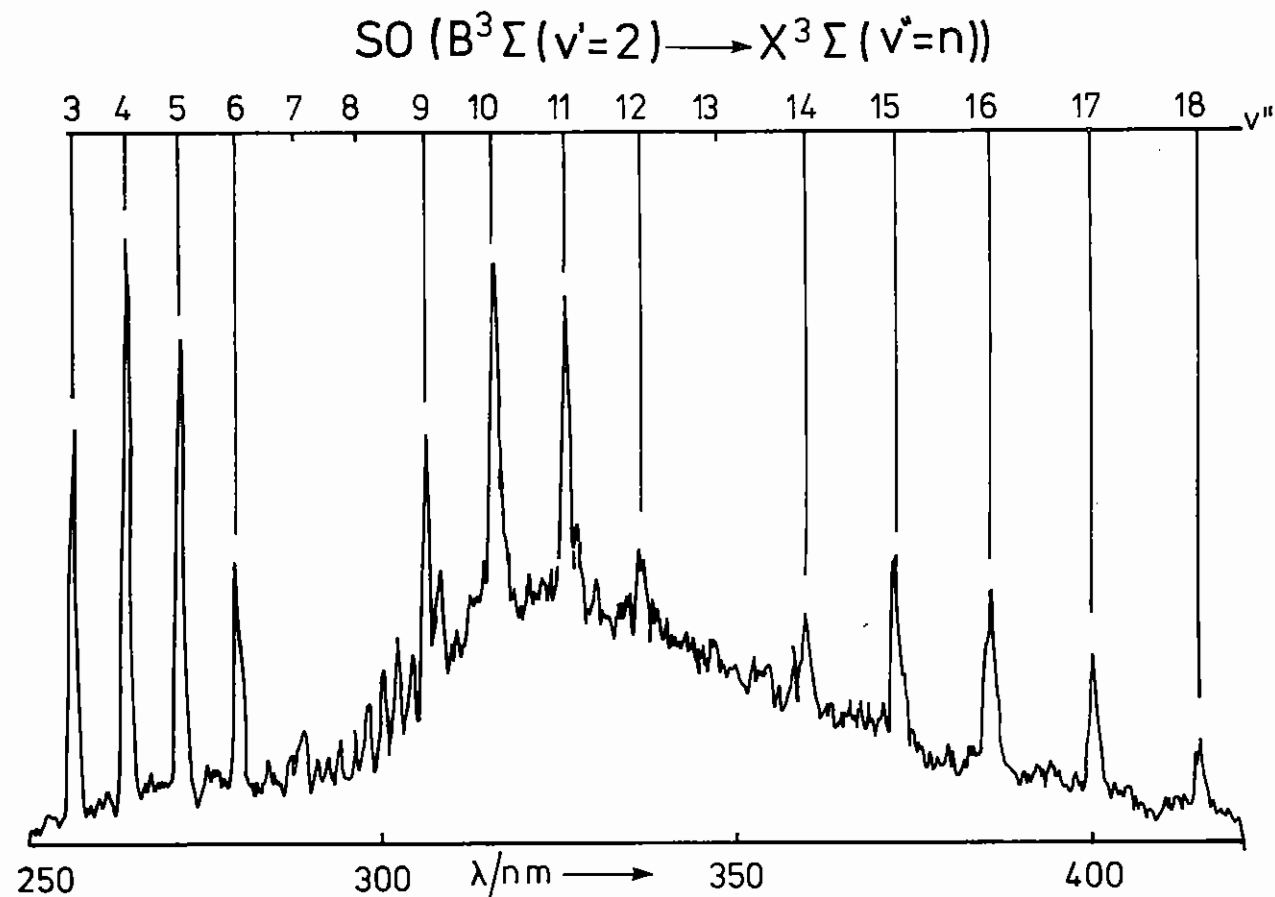


Fig B1.2 Fluorescence observed following excitation of SO_2 ($P = 1.3 \text{ kN m}^{-2}$) at 248 nm with the focused output of a KrF laser. The sharp $\text{SO}(\text{B}^3\Sigma (v' = 2) \rightarrow \text{X}^3\Sigma (v'' = n))$ fluorescence is superimposed on the broader fluorescence from $\text{SO}_2 (\text{B}^1\text{B}_1 \rightarrow \text{X}^1\text{A}_1)$. The spectrum has not been corrected for the response of the detection system.

average power of the model FL 2002 dye laser which has replaced the unsuccessful FL 2001. Replacement of the unsuccessful tetrode switches by specially selected thyratrons in the coupled lasers forming the line-narrowed excimer system led to a large improvement in reliability of this system. Unfortunately, these thyratrons performed acceptably for only ~ 3 months, at which point the 'jitter' became so large as to prevent any line-narrowing. Prior to this, the jitter slowly increased with time, preventing consistent pump/probe experiments with time delays less than ~ 80 ns.

We have been successful in attenuating the KrF laser pulses with Schott glass filters so that a silica flow tube may be used to holed the sample instead of the rapid flow liquid jets described in our last report. This minimizes sample contact with air, which can be important in some experiments, and allows us to reduce the amount of sample needed for a flow experiment from ~ 10 ml to less than 1 ml. Much slower flow rates can also be used, thus extending the residence time of the sample in the "observation zone" and allowing longer delay times between laser pulses to be used if desired.

B1.2.3 Results and Discussion

A diverse set of compounds have been investigated in this reporting period. Examples of TR^3 and resonance Raman spectra of some of these compounds will be given here.

B1.2.3(a) TR^3 Investigations

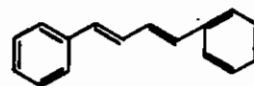
B1.2.3.(a)(i) Anthraquinone-2,6-disulphonate and related anthraquinones

This work was done in collaboration with the group from the Royal Institution and will be described in Section B1.3 of this report.

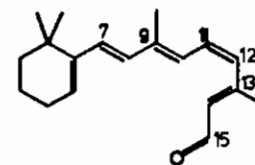
B1.2.3(a)(ii) All-trans-1,4-diphenyl-1,3-butadiene

All-trans-1,4-diphenyl-1,3-butadiene (DPB) is a useful model compound for

the retinal molecule which forms part of the visual pigment rhodopsin. When illuminated in the eye,



all-trans-DPB



11-cis-Retinal

retinal, which is 11-cis in its "dark" form, undergoes a photoinduced isomerization about carbons 11 and 12 to give the all-trans form. In our experiments, all-trans-DPB undergoes a similar excited state isomerization to cis, trans-DPB. We have studied DPB in its excited triplet state. This is produced by efficient energy transfer from biphenyl, triplet biphenyl being readily produced by pulsed 248 nm irradiation of the ground state.

We have obtained the enhanced Raman spectrum of ground state all-trans-DPB in the preresonance region and the resonance Raman spectrum of triplet DPB. The ground state spectrum was obtained using only the dye laser at 396 nm (6 mJ/pulse), while the triplet state spectrum was obtained by pumping the sample and biphenyl sensitizer at 248 nm (10 mJ/pulse) and probing at 396 nm (6 mJ/pulse). Various time delays between the pulses were used. Figure B1.3 shows the TR^3 spectrum of triplet DPB.

This work has been submitted for publication.

B1.2.3(a)(iii) Short-lived alkyl peroxy free radicals and alkyl nitroxyl radicals.

308 nm laser radiation has been used to generate unstable alkyl peroxy free radicals in solution, with line-narrowed, time-delayed 248 nm probe laser radiation being used to excite their resonance Raman spectra. These experiments have proved particularly difficult and, although spectra have been obtained, they are not yet complete.

The alkyl nitroxyl free radicals are more stable than the peroxy analogues and have served us as useful model compounds in this study. We have worked

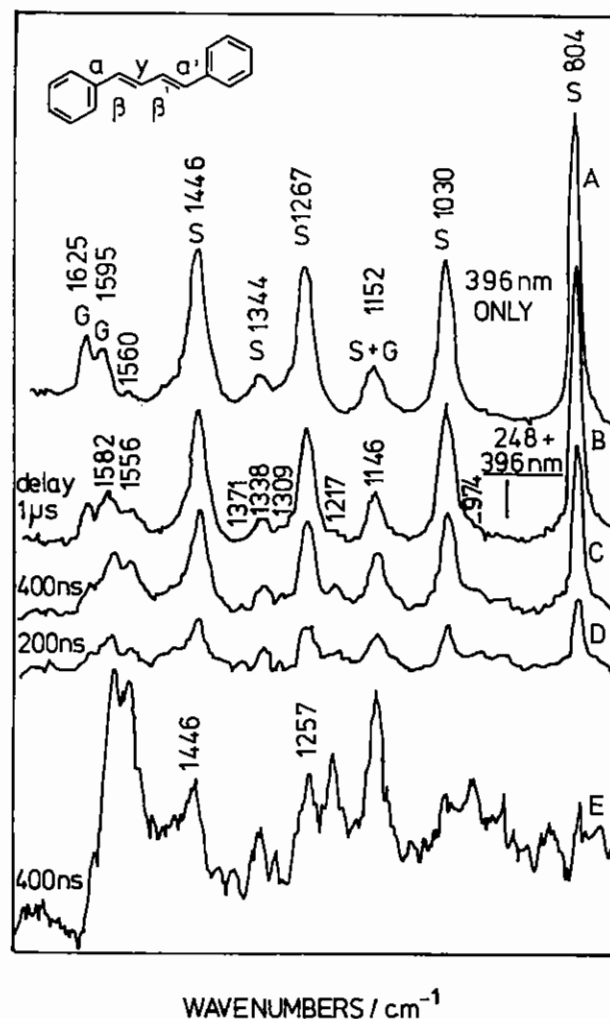


Fig. B1.3 Raman spectra of a solution of 6.10^{-4} M DPB + 8.10^{-4} M BP in cyclohexane, N_2 -saturated. A): probe only; B)-D): pump + probe; B): $1\mu\text{s}$; C): 400ns ; D) 200ns time delay between pump and probe; E): as C) after subtraction of solvent and ground state bands. Intensity scales are arbitrary.

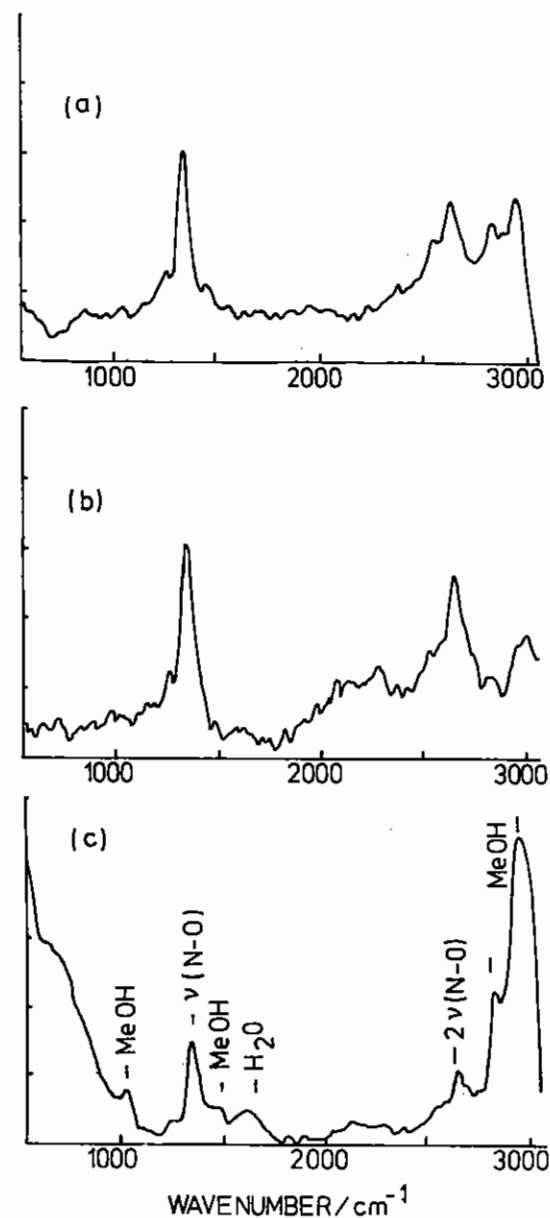


Fig. B1.4 248 nm-excited resonance Raman spectra of (a) DTBNO (5.0×10^{-3} M), (b) TEMPO (3.2×10^{-3} M), both with solvent bands subtracted, and (c) TEMPO (as in (b) but without solvent subtraction).

with dilute solutions in methanol/water of di-*t*-butylnitroxyl (DTBNO), 2,2,6,6-tetramethylpiperidinoxyl (TEMPO) and 4-hydroxy-2,2,6,6-tetramethylpiperidinoxyl (TEMPOL), using 248 nm radiation to excite their resonance Raman spectra. The N-O bond stretching mode has been found at $\sim 1340\text{ cm}^{-1}$ and a number of overtones also have been measured, enabling us to calculate the anharmonicity of the vibrations. Some of these spectra are shown in Figure B1.4. A full account of this work is being prepared for publication.

B1.2.3 (b)(i) Resonance Raman Studies in the ultraviolet (UVRR)
Nicotinamide adenine dinucleotide (NAD^+/NADH).

Resonance Raman spectroscopy has recently been used to probe the active sites of enzymes. These studies all have involved the binding of substrates to form coloured complexes which were then amenable to study with visible lasers. From the biological point of view, the most satisfactory state of affairs would be to use a natural reporter molecule, preferably one which participates in the enzyme reaction. NAD^+ and its reduced form, NADH, are coenzymes for a large number of enzyme-catalysed redox reactions, but have not been exploited as reporter molecules since they are not coloured. NADH does have an absorption at 340 nm, but attempts to excite resonance Raman spectra in this absorption band lead to strong fluorescence; this absorption is not present for NAD^+ . Both NAD^+ and NADH, however, do have a strong absorption at 260 nm. We have obtained resonance Raman spectra of NADH using the line-narrowed 248 nm KrF excimer laser and thus have avoided the fluorescence problems. See Figure B1.5. We plan to use NAD^+/NADH to investigate a number of enzyme reactions, such as that catalyzed by liver alcohol dehydrogenase.

B1.2.3(b)(ii) Papain and Model Compounds

This work was done in collaboration with the group from Birmingham University and is described in Section B2.2 of this Report.

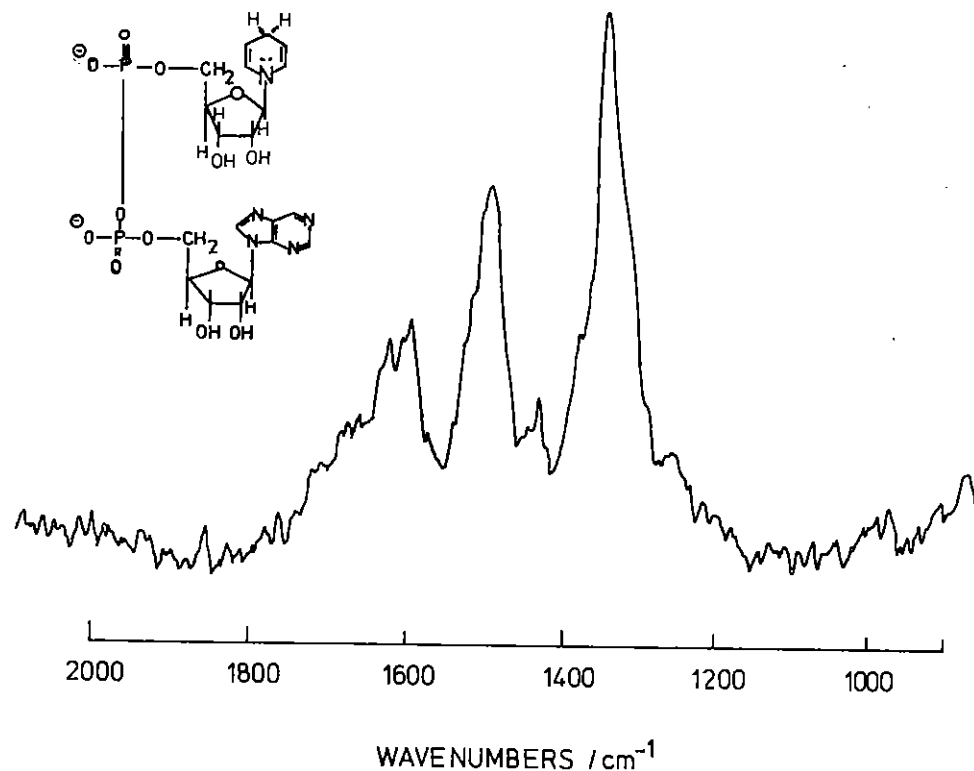


Fig. B1.5 Resonance Raman spectrum of NADH ($5 \times 10^{-4}\text{ M}$) excited at 248 nm.

B1.3 Time Resolved Resonance Raman Spectroscopy of Anthraquinone

D Phillips, J Moore, G Atkinson (Royal Institution), R E Hester, and P Killough (York)

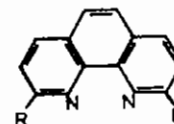
A series of sulphonated anthraquinone derivatives has been studied by time-resolved resonance Raman spectroscopy. A nanosecond pulse from a XeF laser pumps the sample into the excited state, and the system is then probed by observation of the scattered light from a second pulse from an excimer-pumped dye laser. By tuning the dye laser into the absorption bands of the transient species produced by photolysis and using suitable time delays between the pump and probe pulses, the resonance-enhanced Raman spectra of the transient species may be observed. Variation of the time delay between the two pulses allows an estimation of the kinetics to be made, and the vibrational bands observed in the Raman spectra to be identified with the electronic absorption bands seen in flash photolysis experiments. Transient species which have been studied include the triplet, a triplet-solvent complex, and the semiquinone radical anion. Analysis of this data is presently being undertaken with the aim of using the vibrational bands to elucidate the structures of these transient species.

B1.4 Time-resolved Luminescence Spectroscopy of Copper (I) Complexes in Solution

S E J Bell and J J McGarvey (Q U B)

B1.4.1 Introduction

Luminescent charge-transfer excited states of metal-centred systems are of practical and theoretical importance in relation to photophysical and electron-and energy-transfer phenomena. Recent examples, first reported by McMillin et al (B1.6), are the metal-to-ligand charge-transfer (MLCT) excited states of copper (I) complexes with heteroaromatic ligands such as phenanthroline and substituted analogues,



R = H (phen)

R = CH₃ (diMe-phen)

which luminesce in fluid solvents at room temperature. The temperature dependence of the emission has been interpreted in terms of two emitting states in Boltzmann equilibrium. The availability at RAL-UVRF of a wavelength-tunable dye laser and an optical multichannel analyzer affords an excellent opportunity for a comprehensive investigation of the kinetics and photophysics of these interesting systems. The aim of the preliminary experiments described here is to populate the MLCT states of complexes of the type Cu(L)₂⁺ (L denotes phen or diMe-phen shown above) using the excimer-pumped dye laser as the excitation source and to study the luminescent decay kinetics by gating the OMA to time-resolve the MLCT emission spectra.

B1.4.2 Experimental

Figure B1.4.1 shows a simplified schematic diagram of the experimental arrangement. The dye laser was operated in a 'free-running' mode (rep. rate 5-10 Hz) and the accumulation of emission spectra was controlled from the OMA console, utilising timing modes (3 + 7) (as described in the diode array operating manual). Appropriate delays in the range (0-200)ns between an emission event and the gating pulse to the OMA head were selected by means of the built-in delay circuitry on the gate pulse generator. As an independent check, the delay between the gating pulse and that from the photodiode monitor of scattered laser radiation was measured on a storage oscilloscope. In view of the period available (7-14th January 1984), studies were confined to two laser excitation wavelengths, 453nm and 530nm. A Chance glass filter (OR.2) was placed at the entrance slit of the spectrometer (600 grooves/mm grating) to minimise laser scatter. Solutions of the complexes in dichloromethane were degassed by purging with argon and all experiments were conducted at ambient temperature.

B1.4.3 Results and Discussion

The luminescence quantum efficiencies for these systems are low (B1.6), making it necessary to accumulate the signals from several hundred laser shots in each experimental run. Figure B1.7. shows examples of the resultant spectra for $\text{Cu}(\text{diMe-phen})_2^+$ and $\text{Cu}(\text{phen})_2^+$ at several delays following excitation at 530nm. Analogous results were obtained with 453nm radiation. The emission spectrum for $\text{Cu}(\text{diMe-phen})_2^+$ is in good agreement with that reported by McMillin. To our knowledge, that shown for $\text{Cu}(\text{phen})_2^+$ is the first reported emission for this complex in solution. (The emission intensity here is approximately an order of magnitude lower than for the diMe-phen complex). A puzzling feature is the sharp peak at 618nm which decays within 40ns. The other, longer wavelength peak, evident also in $\text{Cu}(\text{diMe-phen})_2^+$ spectra, is attributable to solvent Raman scattering. Studies at other excitation wavelengths and over a wide temperature range will be necessary to establish the origin of the $\text{Cu}(\text{phen})_2^+$ emission.

In the spectra for $\text{Cu}(\text{diMe-phen})_2^+$, the absence of any shift with time in the emission maxima indicates that emission occurs from one level under room temperature conditions, effectively in agreement with McMillin's conclusion. The relaxation times for decay of emission (computed from log plots of spectral peak heights vs. time*) lie in the range (50-90)ns, dependent upon the concentration of parent complex. A detailed analysis of the results (in combination with other kinetic data derived from MLCT excited state absorption studies at Belfast) suggests that the overall photophysical picture is complicated, with perhaps three excited states being involved. Further studies of the time-resolved luminescence at other excitation wavelengths will be necessary to substantiate this conclusion. These trial experiments nevertheless demonstrate the suitability of the high repetition rate dye laser in combination with the gated OMA for photophysical investigation of weakly luminescent systems.

*The validity of this procedure was checked by time-resolving the emission from the MLCT state of the complex $\text{Ru}(\text{bipyridyl})_3^{2+}$ in water. The computed radiative decay time, 500ns, is in very good agreement with literature data.

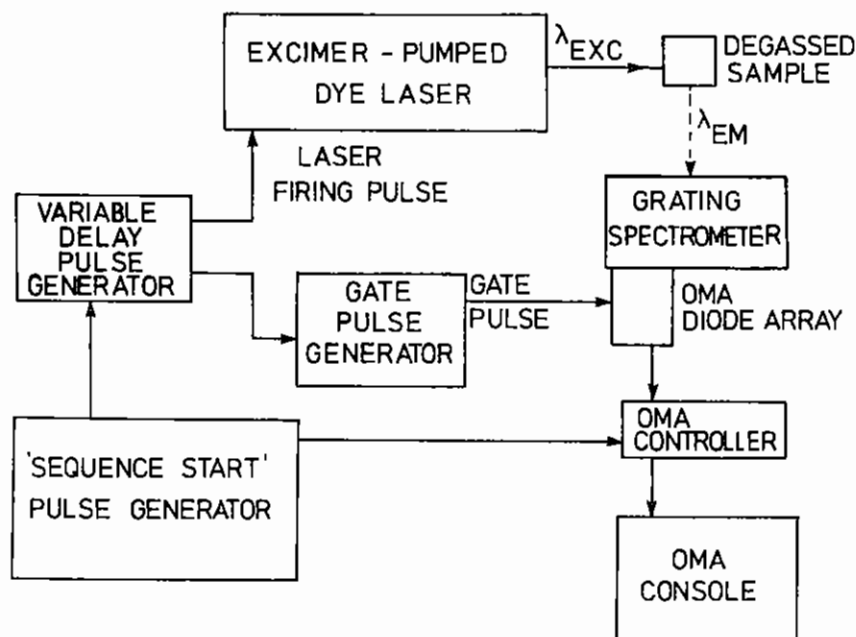


Fig. B1.6

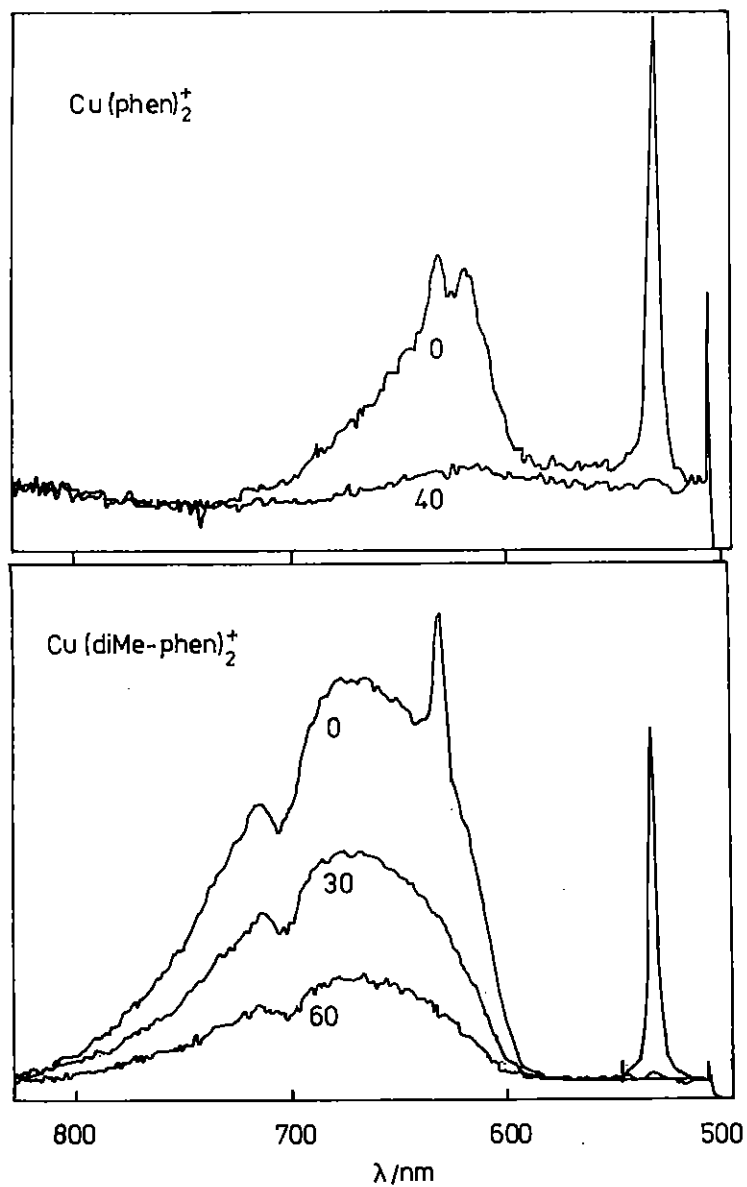


Fig. B1.7

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B2 PHOTOBIOLOGY

B2.1 Time Resolved Fluorescence Polarisation Studies of N-3-Pyrene Maleimide Bound to Bovine Serum Albumin

A D B Waldman, J J Holbrook (Bristol), J Szechi and M H Key (RAL)

B2.1.1 Introduction

N-3 pyrene maleimide is known to form highly fluorescent adducts when reacted with thiol groups and has been used extensively as a fluorescent label in proteins (B2.1) (B2.2).

In this study, conjugates of this dye with Bovine serum albumin have been used as a test system to assess the feasibility of using lasers and streak cameras for measuring the time-resolved fluorescence anisotropy of intrinsic and extrinsic protein fluorophores in the 10^{-10} - 10^{-7} s time range. Such measurements will enable rapid motions of individual residues and domains within globular proteins to be monitored.

B2.1.2 Materials and methods

Pyrene BSA conjugates were prepared from Bovine Serum albumin (Sigma) and n-3-pyrene maleimide (Molecular Probes Inc.) by the method described by Lux and Gerard (B2.3) and dialysed extensively against 0.1M phosphate buffer pH7.5. Samples were all 6.8 μ M with respect to dye concentration (calculated from A_{345nm} assuming an absorption coefficient of $4 \times 10^4 M^{-1} cm^{-1}$ (B2.4). This gave an absorbance of $0.2 cm^{-1}$ at the excitation wavelength (337nm).

Exciting light was provided by a nitrogen-filled Lambda-Physik EMG 201 laser utilising the nitrogen line at 337nm. The intensity of emitted light was resolved in time using a Hamamatsu C-979 streak camera and Cl330 temporal analyser. This was triggered by a fast photodiode and an optical delay line

was employed to give sufficient pre-trigger. Triggering delay was varied by the use of calibrated lengths of coaxial cable.

Vertical polarisation of the laser light was achieved using a Glan-Taylor prism and the emitted light was resolved into horizontal and vertical components by the use of split polarising filters mounted on the inlet window of the streak camera. The horizontal and vertical components of the fluorescence were monitored simultaneously as two streaks on the Vidicon tube and stored digitally in two channels of the temporal analyser. The limited range of streak camera sweep times (maximum of 10ns) meant that it was necessary to concatenate several streaks in order to observe the fluorescence decay over a sufficiently long time to allow lifetime measurements to be made. This, however, was complicated by apparent inconsistencies in the calibration of the full screen width setting on the temporal analyser and the very limited dynamic range of the detection system.

B2.1.3 Results and Discussion

Total fluorescence intensity measurements showed a decay with a lifetime in the range 5-10ns (Figure B2.1). The accuracy with which this could be measured was very limited due to the long duration of the laser pulse (7-10ns FWHM) relative to the fluorescence lifetime. This lifetime corresponds to the shorter lived fluorescence decay component reported by Lux and Gerard (B2.3) which was assigned to n-3-pyrene maleimide covalently bound to BSA via the protein thiol. The longer (50ns) lifetime decay, which these authors ascribed to the presence of dye molecules adsorbed onto the surface of the protein, was not measured due to the narrow range of sweep times and poor dynamic range characteristics of the detection equipment in the configuration used. Anisotropy was clearly detectable in the emitted fluorescence signal (Figure B2.2). The magnitude of the anisotropy changed little in the time range measured which is also consistent with previous studies which have shown rotational correlation times (τ_c) for serum albumin to be in the region of 40ns (B2.5). Again, the poor signal to noise ratio and limited dynamic range of the detection equipment prevented reliable determination of τ_c .

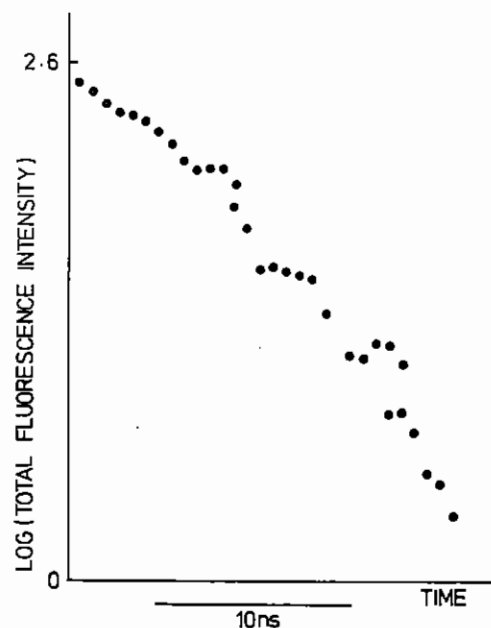


Fig. B2.1 Decay of total fluorescence (calculated as 2H+V) with time. Each data point was calculated from the average of 16 consecutive data points in time. The data shown was taken from three concatenated single-shot experiments, each of 10ns sweep width

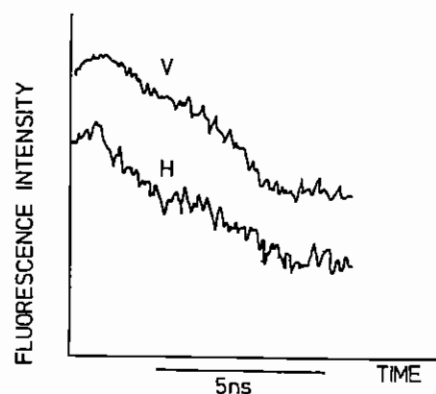


Fig. B2.2 Decay of the horizontally (H) and vertically (V) polarised components of the fluorescence signal with time. Data was obtained from a single-shot experiment.

These preliminary results demonstrate the potential for using laser and streak camera systems for studying time-resolved fluorescence polarisation in proteins. However, accurate measurement of lifetimes and correlation times in the 10^{-10} - 10^{-7} s time range would require laser pulse of short duration (~ 100 ps), a rapidly retriggerable, low-jitter streak camera with larger range of sweep times (50ps-100ns) and a sensitive detection device with wide dynamic range characteristics. Additionally, a continuously tuneable excitation source and some means of monochromating the emitted light are desirable since this allows the contributions to the intrinsic fluorescence of proteins which contain multiple fluorescent residues to be resolved.

B.2.2 Raman spectroscopy of biological molecules

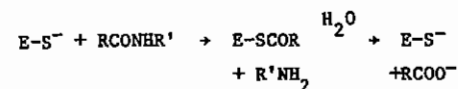
Some preliminary studies of enzyme catalytic centres and mechanisms using ultra-violet resonance Raman spectroscopy

C W Wharton, P J Tonge (Birmingham), R E Hester, P M Killough (York)

B.2.2.1 Introduction

The formal covalent bond exchange processes of enzyme-catalysis are now quite well understood and further progress in our understanding relies upon description of the weaker (H-bonding, steric, hydrophobic) forces that modulate covalent forces during catalysis.

The industrially-important protein-hydrolysing enzyme papain effects catalysis according to the scheme:



The thiolate anion of the essential cysteine residue absorbs in the

ultra-violet with a λ_{max} near 232 nm and has been proposed to exist as a component of a zwitterion with an imidazolium cation of a neighbouring Histidine residue.

The thiol-ester acyl-enzyme also absorbs ultra-violet light with a λ_{max} of 235-232 nm.

Accordingly both the thiolate anion and the acyl-enzyme are amenable to excitation by 248 nm laser irradiation and hence potentially to observation of their resonance Raman spectra. The apparatus used has been described in the 1983 Laser Division Annual Report (Chapter 3, 3.4.2).

B2.2.2(a) Model Compounds

Ethyl thiolacetate has been studied as a simple model of the acyl-enzyme. That the spectrum is resonance-enhanced in dilute aqueous solution is shown in Fig. B2.3.

Thiolate anions (e.g. mercaptoethanol) have been examined in dilute aqueous solution but, as yet, resonance-enhancement of ν (C-S) has not been observed.

B2.2.2(b) Papain

An important finding is that there is little fluorescence interference when the enzyme is irradiated at 248 nm. Also the enzyme retains some 70% of its catalytic activity after exposure to a single laser pulse. All experiments have been conducted using a capillary flow system that, on a bulk flow basis, results in exposure of the enzyme to a single laser pulse. Resonance enhancement of ν (S-S) at 519 cm^{-1} has been observed.

B2.2.2(c) Acyl-papain

A continuous fast flow mixing apparatus has been constructed that allows preparation of the metastable ($t_{1/2} \sim 0.1\text{s}$) acyl-enzyme N-acetyl-L-Phe-Gly-papain using the methyl ester as substrate. Fluorescence

interference is again found not to be serious and resonance-enhancement of ν (C-S) of the thiol ester linkage has been observed at 621 cm^{-1} .

B.2.2.3 Improvements to the Raman apparatus

Based on the results of the preliminary experiments described above, we have identified a number of improvements to our experimental system that are expected to lead to an improved S/N ratio in our presently rather noisy spectra. These are: (1) use of precision-drawn suprasil capillary cells; (2) use of back-scatter collection optics; (3) use of a translatable capillary cell in enzyme experiments to prevent local obscuration of the laser beam.

B.2.2.4 Summary and conclusions

In a short experimental program we have demonstrated that systems described above are amenable to study by resonance Raman spectroscopy using line-narrowed 248 nm radiation. Of particular importance are the facts that observation of the spectra is possible at appropriate (mM) concentrations in water and that fluorescence interference (perhaps rather surprisingly) seems not to represent a problem. A variety of improvements that will be made to the experimental system have been identified.

B2.3 GENERATION AND DETECTION OF SINGLET OXYGEN IN LIPOSOMAL LIPID BILAYERS

A Garner and M J O'Connell (Brunel)

The XeF excimer laser, operating at 351nm and approximately 60mJ, was used to excite anthracene embedded in lipid bilayers that were themselves suspended in aqueous solution as liposomes. Singlet oxygen produced by energy transfer from triplet anthracene to ground state molecular oxygen was detected by observing the formation and decay of the triplet state of β -carotene, itself produced by energy transfer from singlet oxygen.

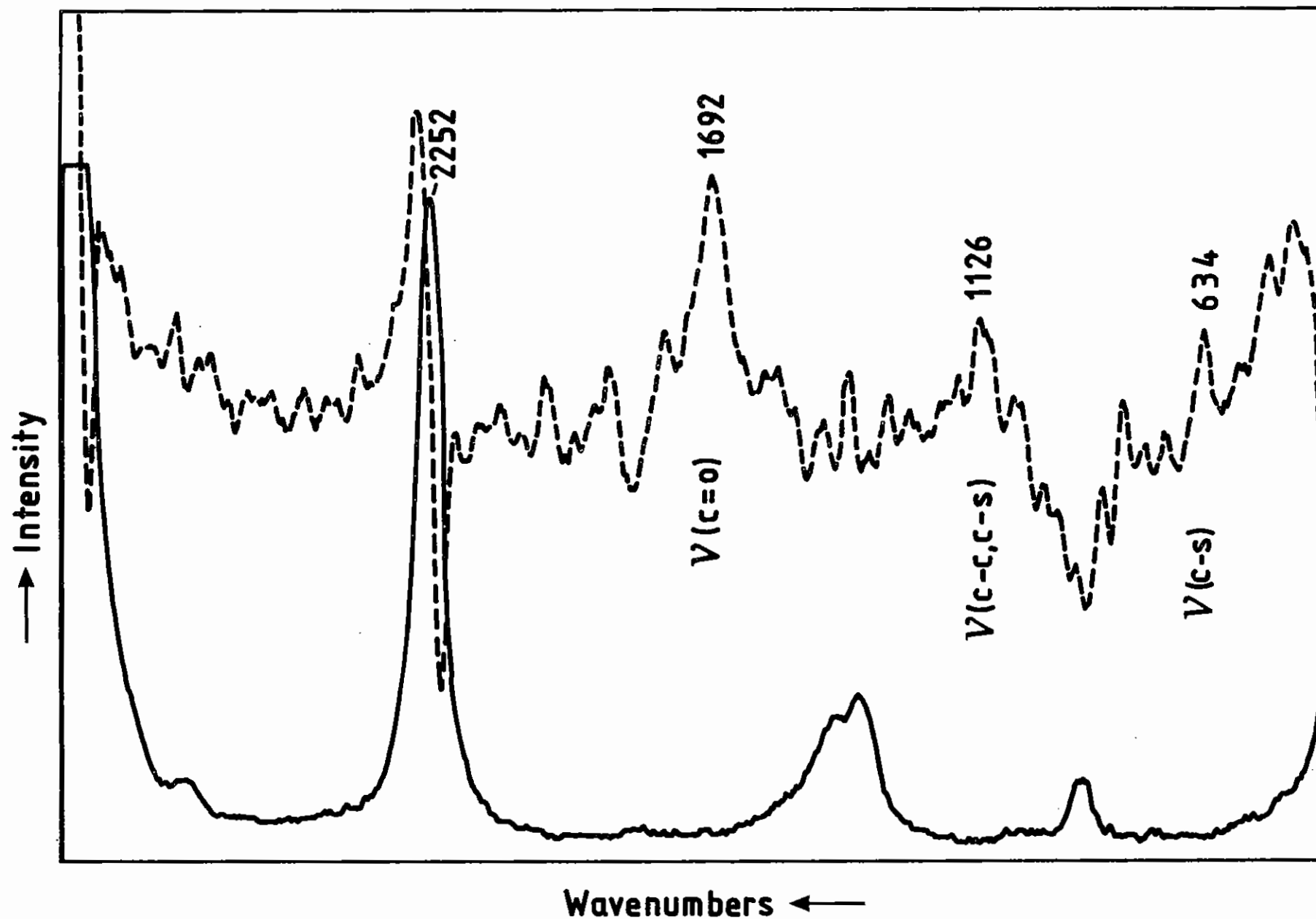


Fig. B2.3

Upper-trace resonance spectrum of ethyl thiolacetate. Lower-trace Raman spectrum of acetonitrile. RR spectrum, 1000 laser pulses, 1 mM solution in 5% acetonitrile/water with solvent subtracted using the 2252 cm⁻¹ band of acetonitrile as null.

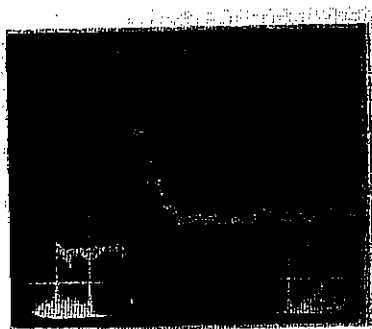


Figure B2.4

The triplet state of β -carotene was monitored by using an argon ion laser, operating at 514nm and 50-100mW, as the probe source. The detector system was exposed to the argon ion laser beam for 2ms during which the excimer laser and the detector recorder system were triggered at appropriate delay intervals. Preliminary results indicate that singlet oxygen in the lipid bilayer has a half-life of about 3 μ s, and is quenched by β -carotene with a rate constant of $10^9 \text{ dm}^3 \text{ mol}^{-1} \text{ s}^{-1}$.

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B3 PHYSICS

B3.1 UV Thomson Scattering on non-Debye discharge plasmas and associated lineshape studies.

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

B3.1.1 UV Thomson scattering

Work has continued on the application of KrF (2480Å) excimer laser radiation to Thomson scattering diagnostics of dense, near non-Debye plasmas (which are optically deep to laser diagnostics in the visible). Initial Rayleigh scattering calibration measurements on Ar and N₂ were reported previously, and used a 150mJ/pulse Lambda Physik EMG 101 KrF laser operating in broadband mode (>5Å). Using the ratio of Rayleigh to Thomson scattering cross-sections, the expected Thomson scattered signal in the ion feature was calculated to be greater than the fluctuation in the plasma emission signal at $n_e = 10^{24} \text{ m}^{-3}$.

Subsequently (May 22 - June 5) an injection locked, line narrowed (<0.1Å) KrF system was used consisting of an EMG 101 grating tuned KrF oscillator and an EMG 201 KrF amplifier. However the low efficiency of this system reduced the pulse energy to 20-30mJ, which was below the level calculated to produce an observable ion feature signal. Hence no signal was observed when Thomson scattering experiments were carried out on the plasma discharge.

However a new line narrowed EMG 150 KrF system producing up to 500mJ/pulse has recently been installed. The output from this system was reduced using apertures and filters to <100 mJ in order to avoid heating of the plasma by more than 5%. A scattered signal level of ~ 6mV was observed above a background pinch emission level of ~ 30mV at peak compression. Using the Rayleigh scattering calibration the scattered signal corresponded to an electron density of $1.3 \pm 0.3 \times 10^{24} \text{ m}^{-3}$. The plasma density at peak compression was diagnosed interferometrically as $1.9 \times 10^{24} \text{ m}^{-3}$ at the centre of the plasma column, which had a FWHM of ~ 2.5 mm. Since the focal

spot size of the laser was also of this order, the slightly lower density given by the Thomson scattering is consistent with averaging over regions of lower density than the pinch core. A search was made for the electron feature between $\Delta\lambda = 15\text{-}29\text{\AA}$ from the laser wavelength ($\lambda_p = 25.5\text{\AA}$ at these conditions), but no such feature was observed. This result was consistent with the ratio of ion-to-electron cross-sections of ~6 at the above conditions ($T_e \sim 2\text{eV}$), giving a predicted 1mV electron feature signal which was less than the fluctuation in the plasma emission ($\sim \pm 2\text{mV}$). Scattering was not observable at higher density conditions due to the poor signal to background ratio.

B3.1.2 VUV Lyman-alpha Lineshape Studies

Successful generation of tunable (1203-1236Å), narrowband (<0.01Å) VUV radiation using a frequency tripled FL2001 dye laser run on butyl-PBD was reported previously. Further experiments in the period June 26 - July 10 demonstrated the feasibility of absorption measurements in the Z-pinch plasma near Lyman-alpha (1209Å). The VUV output from the krypton tripling cell was collimated by a MgF₂ lens, and after passing through the plasma the coherent VUV radiation was incident on the narrow (several micron) slits of a 1m VUV monochromator. The monochromator and its solar blind photomultiplier detector were used to discriminate against the background pinch emission and the untripled fundamental radiation.

Subsequent installation of a more powerful FL2002 dye laser enabled an increase in the coherent VUV output to a level estimated to be $>10^8$ photons/pulse (scheduling period November 20 - December 10 1983). During this period the first Lyman-alpha absorption lineshapes (to our knowledge) were obtained in a dense plasma. The line wing profiles are shown in Figures B3.1-B3.3, the line centre opacity being too great to measure the complete lineshape. The limit to the observable opacity was set by the fluctuation in the pinch emission which exceeded the transmitted coherent VUV signal for opacities >2.5. The large uncertainty in the measured opacities is due to shot-to-shot variation in the intensity of the coherent VUV radiation, the level of which was measured over 20 shots before and after each transmission measurement. It was not possible to monitor the VUV

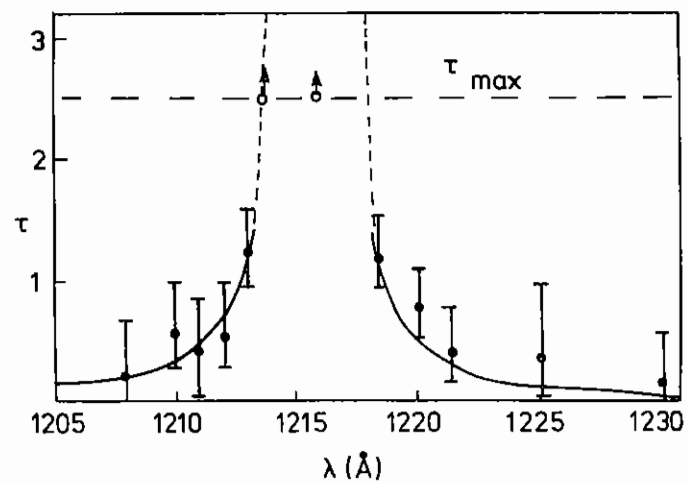


Fig. B3.1 Lyman- α opacity at $n_e = 1.9 \times 10^{24} \text{ m}^{-3}$, $T_e = 2 \text{ eV}$, in a 100:1, Ar:H₂ mixture. Fitted Lorentzian uses a HWHM = 0.1 Å (from Fig. B3.2 giving $\tau_0 \sim 1000$). Dashed straight line gives maximum measurable opacity; open circles indicate a greater opacity.

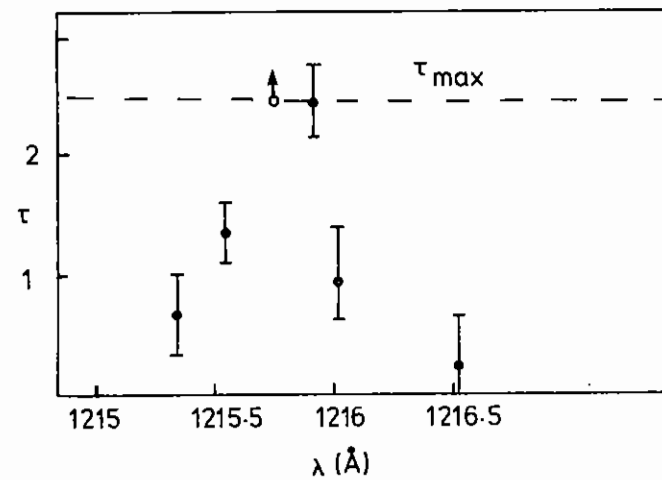


Fig. B3.2 Opacity measurements for the same plasma conditions as Fig. B3.1, but with only a trace impurity of hydrogen in argon.

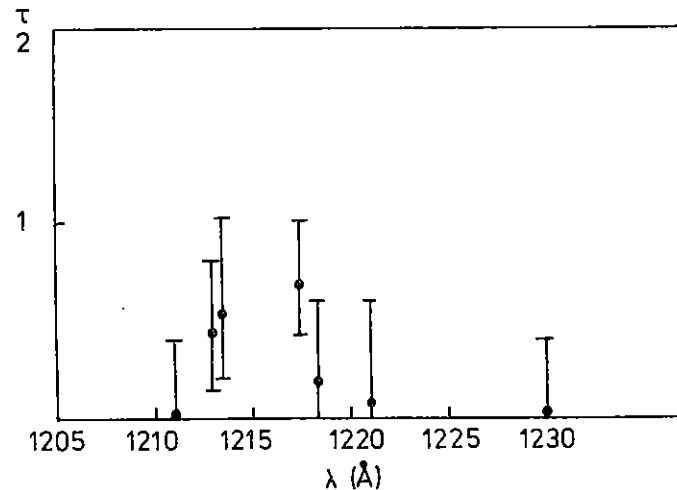


Fig. B3.3 Opacity measurements at $n_e = 4.6 \times 10^{24} \text{ m}^{-3}$
 $T_e \sim 2.5 \text{ eV}$ using 100:1 Ar/H₂ mixture.
 The opacity using only a trace impurity of hydrogen at
 1215.67 was $\tau \sim 0.4$.

intensity of each shot prior to transmission through the plasma, because a small monitor monochromator constructed for this purpose did not produce a signal level which correlated with the level from the 1m VUV monochromator. It was thought that this was due to sampling of different regions of the VUV beam by the two systems, combined with shot-to-shot variation in the spatial intensity profile of the VUV beam.

Measurements of the Lyman-alpha opacity were obtained closer to line centre by using pure (H₂ < 1ppm) argon as the filling gas (impurity H₂ may also have been present in the Z-pinch vessel which was pumped down to $<10^{-2}$ torr before each shot). Figure B3.2 shows the Lyman-alpha absorption profile for the same filling pressure (5 torr) and plasma conditions as Figure B3.1, but with only a trace of H₂. The profile indicates a maximum HWHM for Lyman-alpha of approximately 0.1Å, which is less than the estimated line width at $n_e = 2 \times 10^{24} \text{ m}^{-3}$ ($\sim 0.4 \text{ Å}$).

This result may indicate that the main Lyman-alpha absorption is taking place outside the region of highest electron density at the Z-pinch core.

The upper limit placed on the Lyman-alpha HWHM from the "pure" argon measurement allowed a Lorentzian line shape to be fitted to the line wing profile in Figure B3.1 (a Gaussian fit to the measured points using a HWHM of 0.1Å yielded a profile which decreased too rapidly in the line wings). The Lorentzian profile fits the blue wing quite satisfactorily, but the measured opacities in the red wing lie consistently above the Lorentzian fit. This result supports the red wing asymmetry observed previously in the same Z-pinch device under the same plasma conditions.

A similar asymmetry can be seen in the Lyman-alpha line wing profile at the higher density plasma conditions shown in Figure B3.3. However the measured opacities are lower at higher density (and higher temperature) consistent with a lower $H_{n=1}$ population and a broader line width. Measurement of the opacity at the nominal Lyman-alpha line centre (1215.67Å) using "pure" argon as the filling gas yielded a value of 0.4, which is also lower than the opacity measured at this wavelength for the lower density condition. Although lack of further scheduled experimental time meant that more

Lyman-alpha lineshapes could not be obtained, the present experiments have demonstrated for the first time the feasibility of high resolution absorption line shape measurements in dense plasmas using coherent radiation in the VUV.

B.3.2 Phase Conjugation Studies

G Davis, C E Webb (Oxford), M C Gower and A F Gibson (RAL)

B.3.2.1 Degenerate four-wave mixing using phase conjugate pump waves

The most commonly used techniques for generating phase conjugate (PC) signals are degenerate four-wave mixing (DFWM) and stimulated Brillouin scattering (SBS). However, these two processes have contrasting requirements for producing high quality wavefront correction. Using DFWM, high quality pump beams are required for good reproduction of the probe beam. On the other hand, while SBS readily conjugates aberrated pump beams, unlike DFWM it is a threshold process and usually requires much higher intensities. Features as small as $\sim 0.6\mu\text{m}$ have been reproduced using a DFWM phase conjugate (B3.1) mirror while we have (see below) reproduced feature sizes of $\leq 10\mu\text{m}$ by SBS.

For counterpropagating pump waves, E_1 and E_2 , and a probe wave, E_3 of the same frequency ω , the induced third-order nonlinear polarisation responsible for phase conjugation by DFWM is given by

$$P_4^{(3)}(r) = \frac{1}{2} \chi^{(3)} e^{i(\omega t + k_3 \cdot r + \phi_1 + \phi_2 - \phi_3)} \epsilon_1 \epsilon_2 \epsilon_3^* + \text{c.c.} \quad (\text{B3.1})$$

for input waves of the form $E_i = \frac{1}{2} \epsilon_i(r_i) e^{i(\omega t - k_i \cdot r + \phi_i)} + \text{c.c.}$ $\chi^{(3)}$ is the third order susceptibility and ϕ is the phase of the wave. From Equation (B3.1) it can be seen that the radiated field, E , has a wavenumber $-k_3$ and

a phase $\phi_4 = \phi_1 + \phi_2 - \phi_3$. Thus for uniphase plane pump waves so that ϕ_1 and ϕ_2 are independent of r , perfect phase conjugation (i.e. $\frac{d\phi_4}{dr} = -\frac{d\phi_3}{dr}$) may be achieved using DFWM. However, it has also been recognised for some

time (B3.2) that the drawback of having to use aberration free pump beams for high fidelity phase conjugation by DFWM might possibly be overcome by using pump beams which are phase conjugates of each other (i.e. $\phi_2 = -\phi_1$).

By using an SBS retroreflecting mirror to generate the phase conjugate of aberrated counterpropagating pump waves, we now report high quality PC signals using a DFWM geometry - since the linewidth of the laser used in our study was approximately twice the frequency shift of the Brillouin medium, the four-wave mixing remains still essentially degenerate. Since our DFWM medium need not be Brillouin active, our experiment is conceptually different to the strong/weak pulse non-threshold SBS experiment reported by Basov and Zubarev (B3.2) and further analysed by Abrams et al. (B3.3).

Our experimental setup, which is shown in Figure B3.4, used a KrF line-narrowed, diffraction-limited oscillator operating at 249 nm. In order to obtain a good quality high intensity beam, the 5 ns FWHM duration pulse from this oscillator was then amplified in a double-pass of amplifier A using an SBS PC hexane mirror. We have previously shown (B3.4, B3.5) that this type of mirror can correct for aberrations induced on the wavefront by the amplifier. A standard DFWM retroreflecting geometry was then used with Rhodamine 6G (Rh6G) diluted in ethanol (10^{-5} M concentration) in the DFWM cell.

In earlier work (B3.6) we found such an arrangement could produce PCR's of up to 300% by a thermally induced third-order susceptibility. However, as can be seen in Figure B3.4, in the present experiments a second SBS mirror was used to provide the counterpropagating pump wave, I_2 , which was the phase conjugate of the forward travelling pump I_1 . In some cases, in order to increase I_2 , a second KrF amplifier (amplifier B in Figure B3.4) was placed between this SBS mirror and the DFWM cell.

In Figure B3.5(a) we show a photograph of the phase conjugate beam obtained from the DFWM set up when a conventional plane 50% reflecting mirror was used to derive the second pump wave. In this case, with the pump waves not being phase conjugates of each other, when a fused silica aberrator was placed between the DFWM cell and this mirror, the quality of the phase

DFWM WITH PHASE CONJUGATE PUMP WAVES

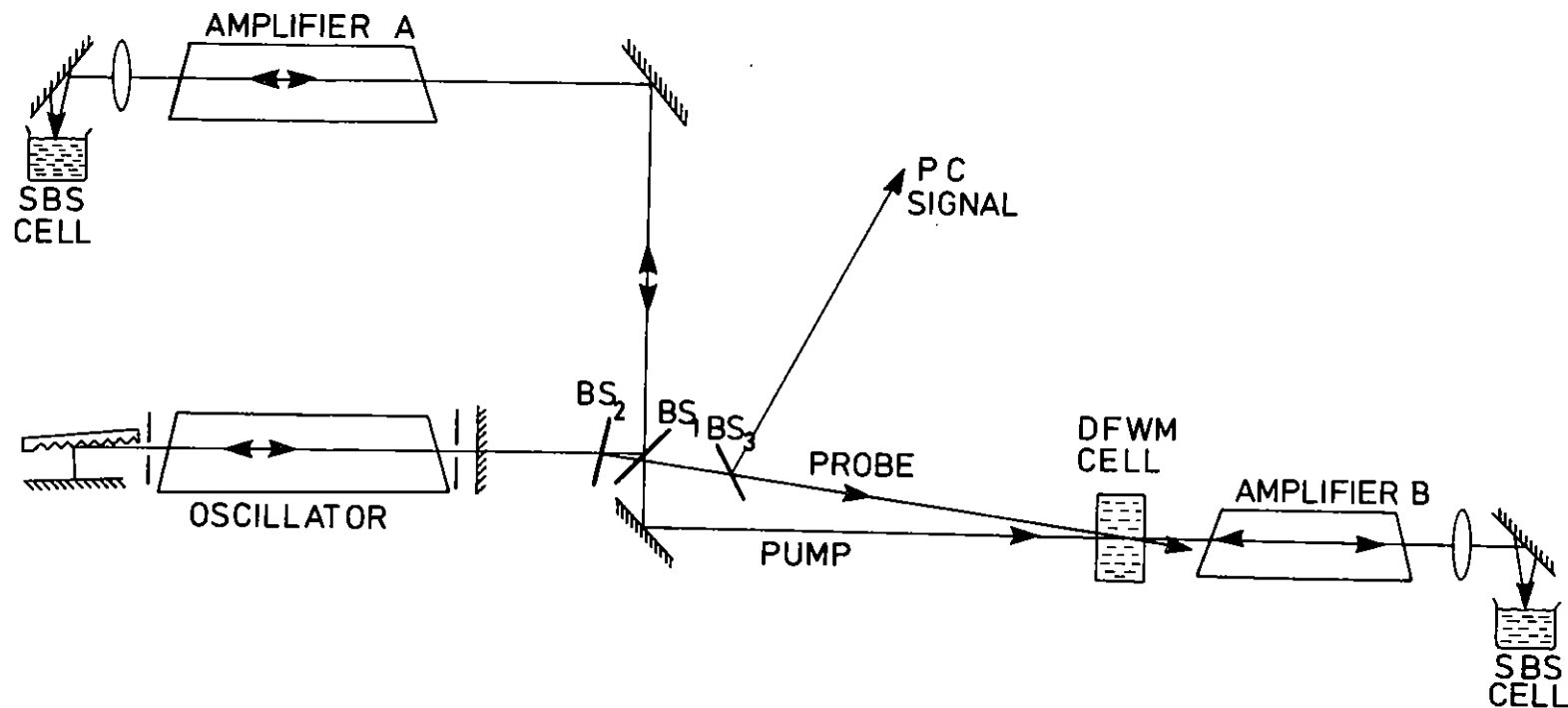


Fig. B3.4 Experimental setup. BS₁, BS₂ and BS₃ are 20%, 50% and 4% reflecting beam splitters, respectively.

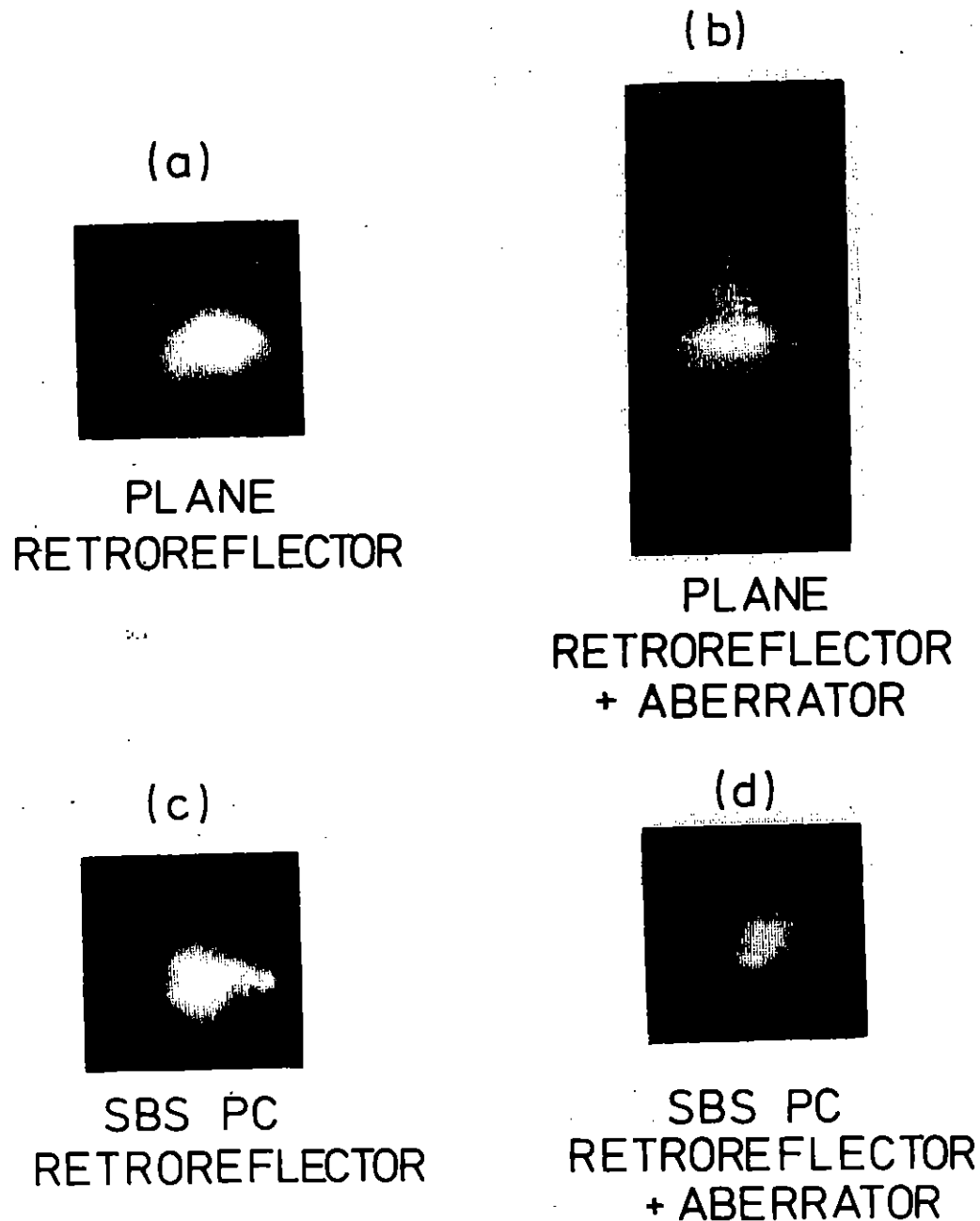
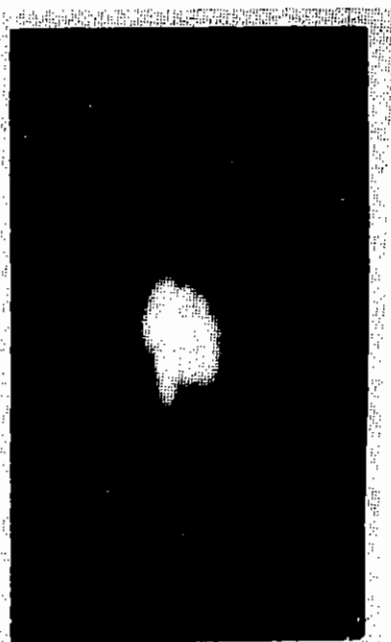


Fig. B3.5

Using Polaroid Type 55 film, photographs of PC beams generated by DFWM. 10^{-5} M Rh6G in ethanol. $I_1/I_3 \sim 4$, $I_1 \sim 12 \text{ MW/cm}^2$.

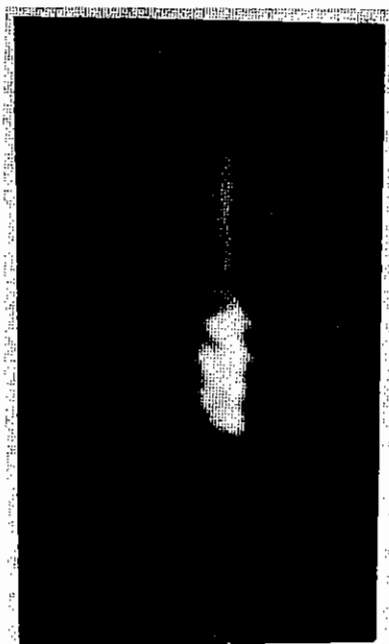
conjugation derived by DFWM becomes very poor - as shown in Figure B3.5(b). On the other hand, as shown in Figures B3.5(c) and B3.5(d), when the plane mirror was replaced by an SBS PC mirror, so that the pump waves were phase conjugates of each other, the quality of the phase conjugate beam derived by the DFWM process remained unaffected when the aberrator was inserted into the path of the second pump wave. Furthermore, when amplifier B, as shown in Figure B3.4, was put into position so that I_2 was increased, as shown in Figure B3.6 similar effects were observed. As we have previously shown, (B3.4, B3.5) such amplifiers can severely aberrate beams travelling through them. This is apparent in Figure B3.6(b) which shows the poor quality phase conjugate signal derived by DFWM when a conventional plane mirror is used behind the amplifier to reflect the pump wave. However, as is apparent in Figure 3.6(c) the quality of this signal is restored when an SBS mirror is used as a retroreflector behind the amplifier to provide phase conjugate pump waves. This clearly demonstrates the utility of using phase conjugate pump waves in DFWM to provide high quality image reconstruction. In Figure B3.7 we show a plot of the PCR versus I_1 obtained from our DFWM setup using both a 50% conventional plane reflector and an SBS PC mirror to derive the second pump wave. A variable pressure NO_2 ($\sigma \sim 4 \times 10^{-20} \text{ cm}^2$ at 249nm) absorption cell placed at the output of the amplifier was used to vary the laser intensity. As can be seen the DFWM reflectivity is approximately the same using either reflector. The maximum observed PCR of $\sim 2\%$ is to be compared with the value of $\sim 11\%$ which we measured previously (B3.7) using three times longer laser pulses of the same intensity and a conventional plane mirror retroreflector. We recall that in the low reflectivity regime ($R \lesssim 1\%$), the PCR derived by DFWM using a thermal mechanism is proportional to the square of the pulse duration (B3.6). For low reflectivities, the DFWM theory for the thermal mechanism also predicts a quadratic dependence of PCR on pump wave intensity. The data shown in Figure B3.4 appears to agree with this prediction (when using an SBS retroreflector once above threshold the reflectivity of this mirror is fairly insensitive to laser intensity (B3.4, B3.5).

(a)



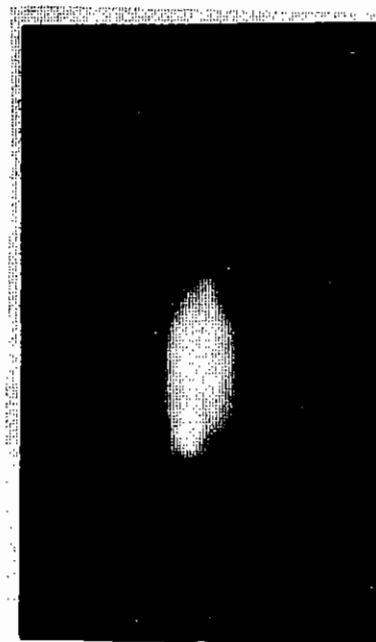
PLANE
RETROREFLECTOR

(b)



PLANE
RETROREFLECTOR
+ AMPLIFIER B

(c)



SBS PC
RETROREFLECTOR
+ AMPLIFIER B

Fig. B3.6 Photographs of PC beams generated by DFWM with amplifier B in place.

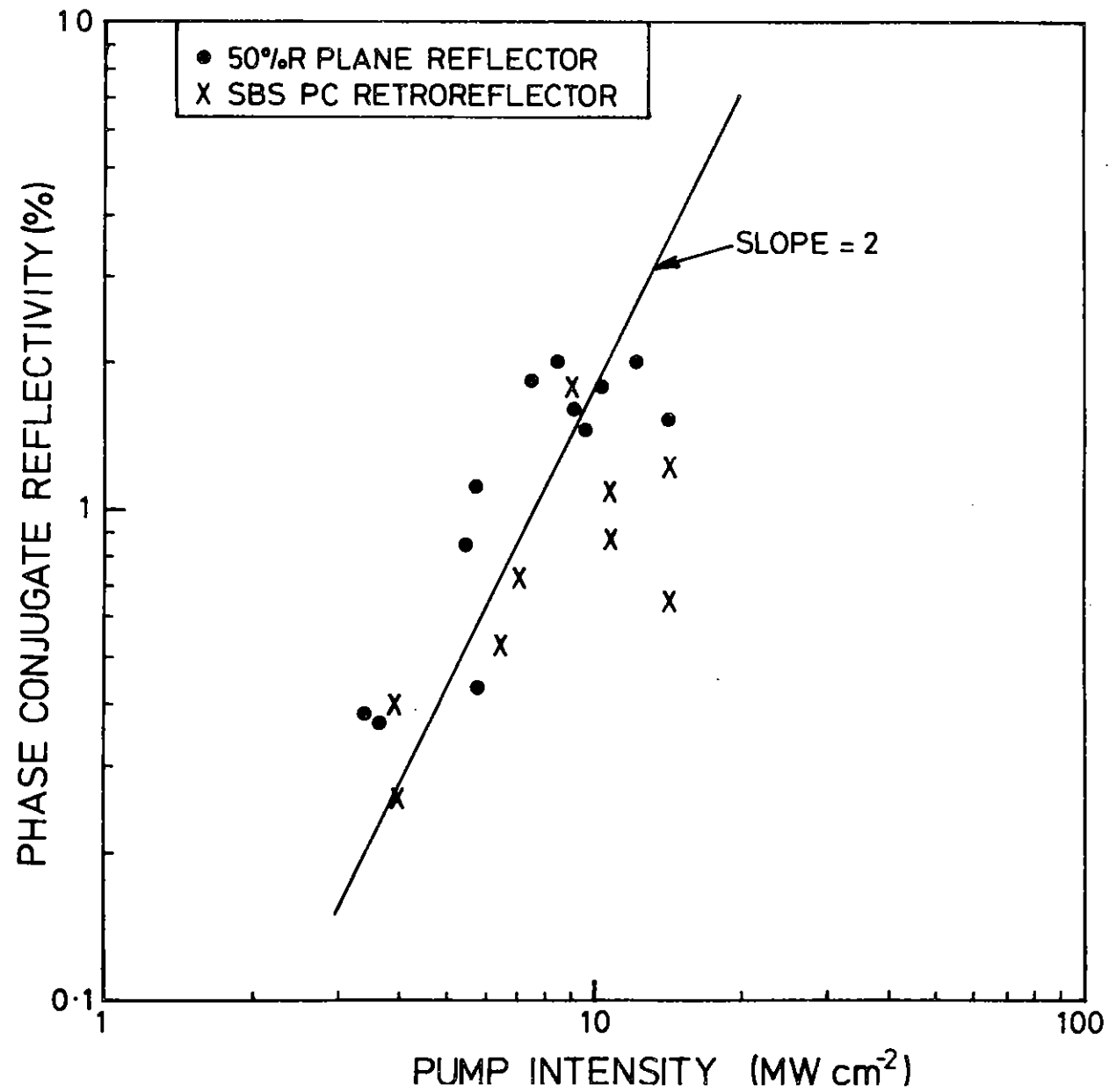


Fig. B3.7 DFWM PCR at 249nm versus pump wave intensity using a 50% reflectivity plane mirror (●) and an SBS PC reflector (X) as retroreflectors. $I_1/I_3 \sim 4$.

In conclusion, using an SBS PC retroreflector to derive the second pump beam, we have demonstrated the utility of using phase conjugate pump waves to obtain good quality phase conjugate signals from a DFWM geometry which has aberrated pump waves. The maximum measured DFWM PCR using this arrangement was ~2% and exhibited approximately a quadratic dependence on pump wave intensity. We are currently performing further experiments to determine the ultimate resolution of this setup.

B3.2.2 Imaging Projection using a PCM

Using an arrangement like that shown in Figure B3.8 it is possible to project images from the beam splitter onto a substrate which is in the same equivalent plane as the (mask) object. Interest in such a technique has thus far been motivated by its potential application to photolithography as used in the production of silicon chips. In this application features as small as 0.5-1 μ m are required to be projected over a 1x1cm² mask area. Conventional lens or mirror systems capable of achieving this are extremely complex and expensive and at present do not work in the ultraviolet (<320nm) spectral region where (to achieve higher resolution) one would like to operate. PCM's are attractive as a potential method of performing cheap, simple high-power UV image projection with a large numerical aperture (NA). In Figure B3.9 we show a setup used by Levenson(B3.1) to project images with ~800 lines/mm resolution at 355nm with a NA ~ 0.19. The PCM was produced by DFWM via thermally induced gratings in Rh6G dissolved in acetone. Although the reflectivity of such a mirror can reach 300% (B3.6), because of the low average power of his laser (~1W), and associated damage problems to the beam splitter cube and the mask at higher powers, the minimum exposure time for the photoresist using this system was ~1min, which is much too long for any practical photolithography system. To circumvent these problems, we are currently testing a system which will employ a KrF laser and a small amplifier between the PCM and the beamsplitter. This system has been designed to produce exposures of the photoresist by a single 20nsec laser pulse. In Figure B3.10 we show photographs of the projected images of 40 and 10 μ m lines using an SBS PCM operating at 249nm. We believe this to be the highest resolution yet demonstrated for SBS PCM image projection. Due to difficulties involved in producing fine line masks, image reduction using

a PCM would be a great bonus for such a device. Although image reduction has been achieved using a four-wave mixing PCM(B3.8) by writing the hologram at one wavelength and reading at a shorter wavelength, such a technique is of limited use, since now the probe wave must be phase matched with E_1 and E_2 . This severely restricts the numerical aperture of such a device. However, it should be possible to construct a practical image reducing PCM having a resolution of ~10⁶ pixels. The main disadvantage of this device is that because of the phase matching requirement, the conjugate wave is no longer counterpropagating to the object wave and so the mirror cannot be used in double-pass arrangements and still achieve aberration compensation. Thus one of the most attractive features of the PCM would be lost. One way imaging through an aberrator using DFWM has also been demonstrated(B3.9), although it is possibly only useful for thin aberrators.

B3.3 Stimulated Anti-Stokes Raman Scattering and Stimulated Collisional-Induced Fluorescence in Atomic Thallium

D G Cunningham, D Denvir, I Duncan and T Morrow (QUB)

Stimulated anti-stokes Raman scattering in inverted media is of current interest as a potentially efficient route for laser frequency upconversion into the VUV (B3.10 - B3.13). The advantages of such up-converters stem from their inherent tunability and from the fact that energy stored slowly in the system may be extracted quickly and efficiently by the application of a pump laser pulse.

However when the photon energy of the pump laser is such that $h(\omega_{23} - \omega_p) < kT$ and the anti-stokes up-converter is in a collisional environment, then the resonance level may be populated by inelastic collisions (see Figure B3.11). These inelastic collisions transfer atoms from the laser induced virtual level to the upper resonance level removing photons from the pump laser and depleting the population inversion employed by the anti-stokes up-converter. Thus in a collisional environment near resonance a drop in the conversion efficiency of the anti-stokes up-converter is to be expected. Under suitable conditions, the collisionally pumped resonance level may be inverted with respect to the lower level of the Raman up-converter and a

IMAGE PROJECTION

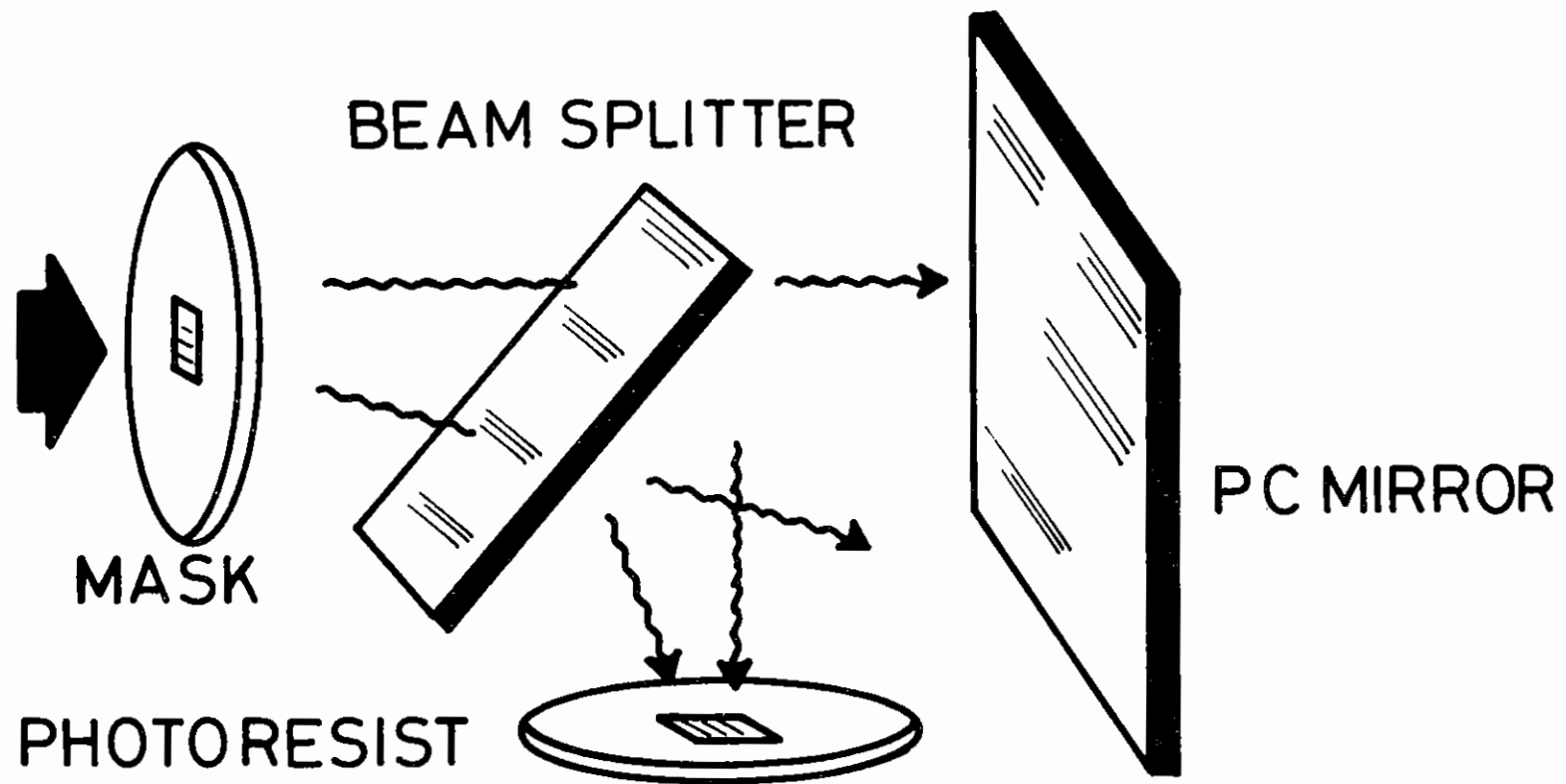


Fig. B3.8 Image projection using a beam splitter and a phase conjugate mirror.

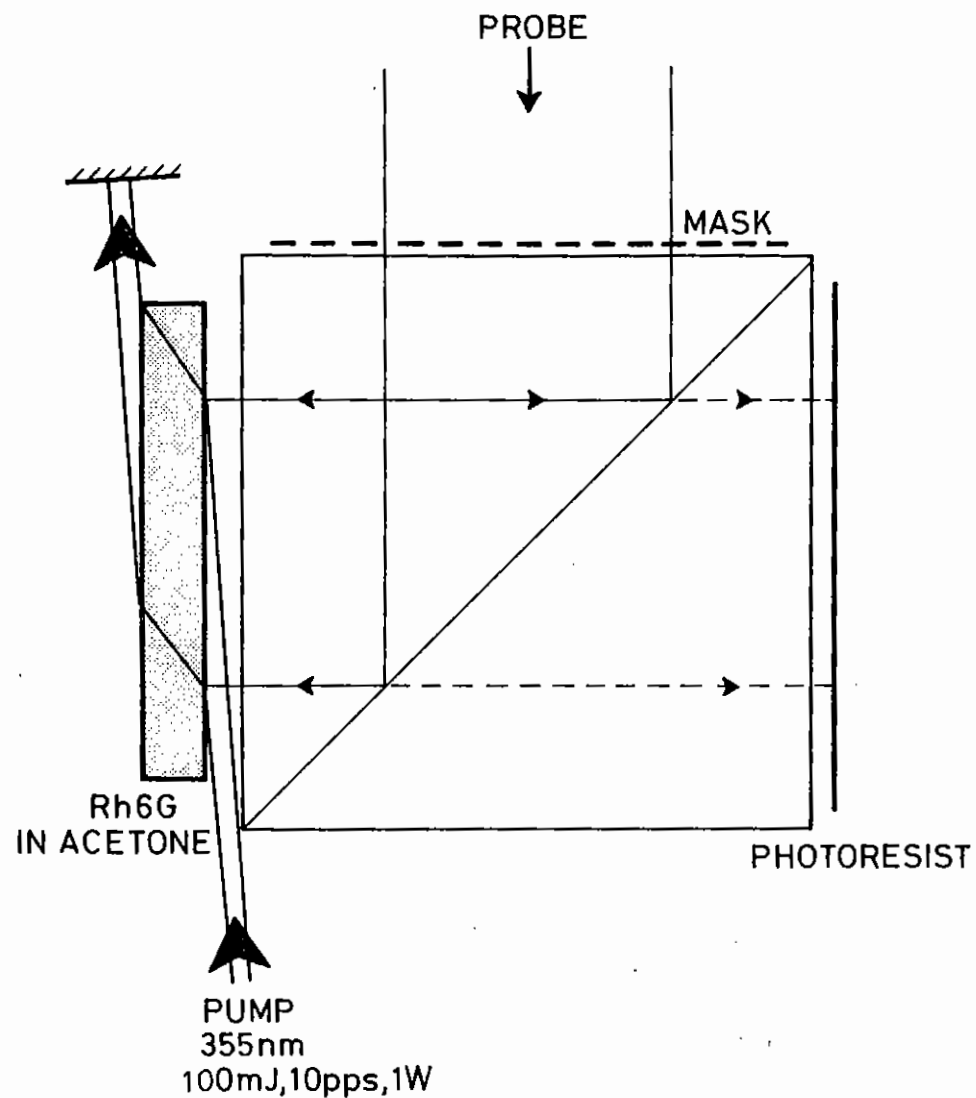
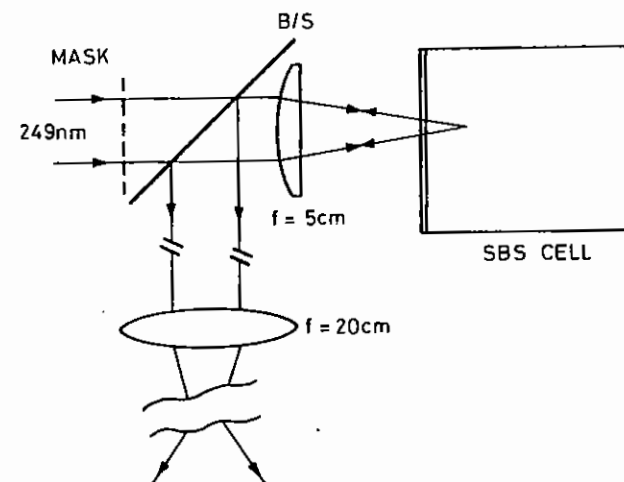


Fig. B3.9 Image projection of a mask onto a photoresist using DFM at 355nm in Rh6G solution in acetone and a beam splitter cube. After Reference B3.1.



40μm LINES AND SPACES



10μm LINES AND SPACES

IMAGE REPRODUCTION BY SBS IN HEXANE AT 249nm

Fig. B3.10 Image projection of 40μm and 10μm line and space masks used SBS in hexane at 249nm.

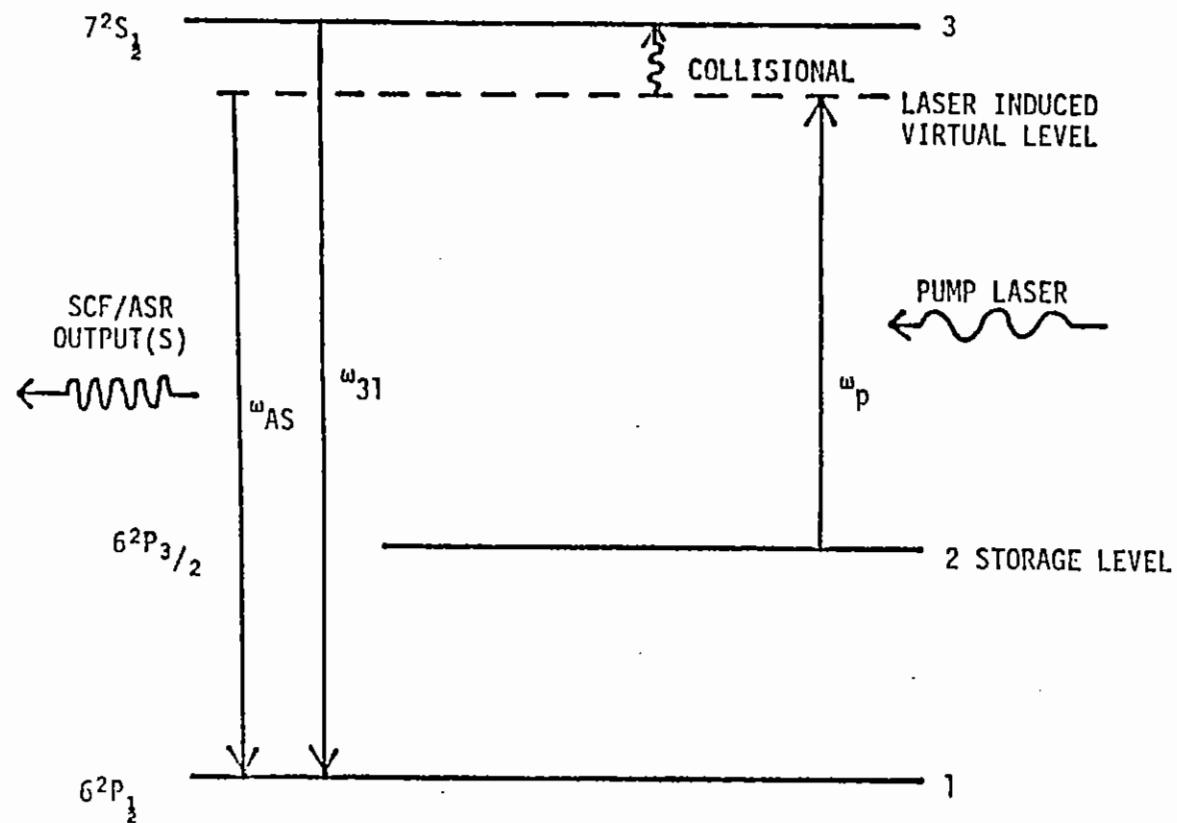


Fig. B3.11 Schematic representation of Anti-stokes Raman and Stimulated collision induced up-converters. (The term symbols for the corresponding levels in Tl are indicated).

stimulated collisional induced laser produced. Critical factors which determine the population mechanism and operating conditions of the SCF up-converter are the collision width γ_c of the intermediate level, $\Gamma_c (= 1/\tau_c)$ where τ_c is the collisional duration, usually of the order of 1 psec) and $\delta (= \omega_{23} - \omega_p)$. If $\delta \leq \gamma_c$ then the SCF can be considered to be amplified spontaneous emission (ASE) due to direct optical pumping (i.e. the pump laser is tuned within the pressure broadened transition), but if $\delta \gg \gamma_c$ the SCF is not due to direct optical pumping. For $\delta \ll \Gamma_c$ both the impact and Markovian approximations are valid, but if $\delta \gtrsim \Gamma_c$ then the SCF up-converter is operating in the non-impact, non-Markovian regime (B3.14). In our experiments $\gamma_c \sim 0.1 \text{ cm}^{-1}$ and $\Gamma_c \sim 10 \text{ cm}^{-1}$ (B3.15). Our experiments, reported below, investigated the stimulated collisional induced fluorescence at 377.6 nm produced in a system of Tl ($6p^2P_{3/2}^0$) atoms. The Tl ($6p^2P_{3/2}^0$) inversion was created by selective photodissociation of TlCl at 193 nm by an ArF laser. i.e. $\text{TlCl} + \lambda (193 \text{ nm}) \rightarrow \text{Tl}^* + \text{Cl}$. The experiments were carried out with excess pressures of argon and the influence of the 'collisional' pumping of the Tl ($7S_{1/2}$) level on both the anti-stokes Raman and stimulated collisional induced outputs was investigated. The SCF generation was observed to extend into the non-impact, non-Markovian regime.

The experimental arrangement is illustrated in Figure B3.12. The 193 nm radiation used to dissociate the TlCl was generated using a commercial ArF excimer laser. Approximately 30 mJ of 193 nm light was focussed into the oven. The ArF laser pulse duration was ~ 10 nsec and was focussed to an area of about $3 \times 10^{-2} \text{ cm}^2$ over the 10 cm hot zone in the cell. The pump laser, tunable around 535 nm, was provided by an excimer pumped dye laser. It had a pulse length of ~ 10 nsec, a line width $\sim 0.2 \text{ cm}^{-1}$, was focussed to an area $\sim 10^{-3} \text{ cm}^2$ over the hot zone and 2.5 mJ of 535 nm radiation entered the cell. The pump dye laser beam was focussed with a separate lens and combined with the ArF laser using a flat with a small aluminized spot, the aluminium coating was protected from possible damage by the Ar F laser by a small shield. In this configuration the pump dye laser beam defined the interaction region. Temporal overlap of the two beams was controlled by a delay generator giving a jitter of < 5 nsec between the outputs of the two

lasers. The TlCl cell was made of pyrex lined with a stainless steel mesh wick and had cold, 'O' ring sealed, unaligned spectroslil 'B' windows. Argon was used as a buffer gas at two pressures, viz 28 and 68 Torr. The TlCl vapour density was monitored using the absorption of the 253.7 nm line of a mercury lamp. The spectral profiles and intensities of the anti-stokes and collisional induced outputs as a function of dye laser wavelength were studied using a 0.5 m spectrometer with an image intensified diode array (OMA) system. Fast diodes were used to monitor the temporal outputs of each laser.

At TlCl vapour pressures $\sim 1 \times 10^{16} \text{ cm}^{-3}$ and with buffer gas pressures of 28 and 68 Torr the stimulated collisional induced laser emission dominated the output of the system. The SCF output energy as a function of pump laser wavelength is shown in Figure B3.13. For each curve the pump laser intensity and Tl($6p^2P_{3/2}^0$) density remained constant only the buffer gas pressure was different. The spectral profile of the SCF output for any particular pump wavelength showed detectable asymmetry, even with the limited resolution available ($\sim 0.1 \text{ nm}$), in contrast to the spectral profile of the ASR output which was symmetric. The asymmetry was particularly marked at large detuning and near threshold. The asymmetric spectral profile of the SCF output could be indicative of interference effects between the SCF and ASR processes or Tl*-Ar quasi-molecule formation (B3.15). At a buffer gas pressure of 28 Torr the tuning curve of the SCF output has a half width $\leq 2 \text{ cm}^{-1}$ and direct optical pumping is important at this pressure. At a buffer gas pressure of 68 Torr the half width of the SCF tuning curve is $\sim 20 \text{ cm}^{-1}$ and so direct optical pumping is not an important contribution for most of this curve (i.e. for $\delta \sim 1 \text{ cm}^{-1}$) but will contribute close to resonance. Since the SCF output at 68 Torr was obtained for $\delta \sim 50 \text{ cm}^{-1}$, we are clearly in the non-impact, non-Markovian regime.

At TlCl vapour pressures between 10^{16} and $3.5 \times 10^{16} \text{ cm}^{-3}$ both the ASR and SCF emissions were detected. As the TlCl vapour pressure was increased the contribution of ASR generation began to dominate the output until at $4 \times 10^{16} \text{ cm}^{-3}$ little SCF emission at 377.6 nm was observed. A tuning curve of the ASR laser (output energy as a function of pump dye laser wavelength) for a buffer gas pressure of 68 Torr and TlCl density of $3.6 \times 10^{16} \text{ cm}^{-3}$ is

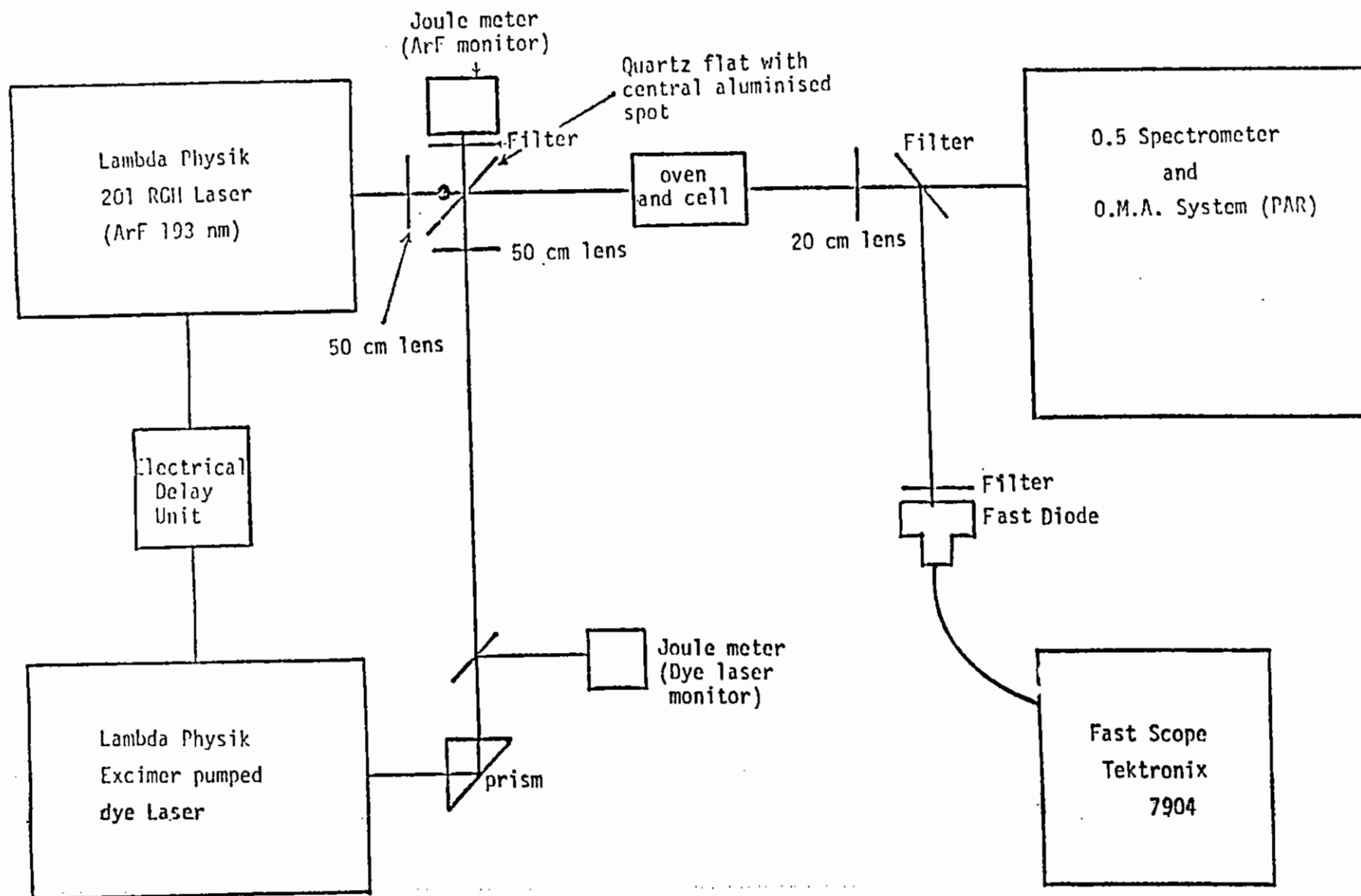


Fig. B3.12 Schematic diagram showing experimental arrangement.

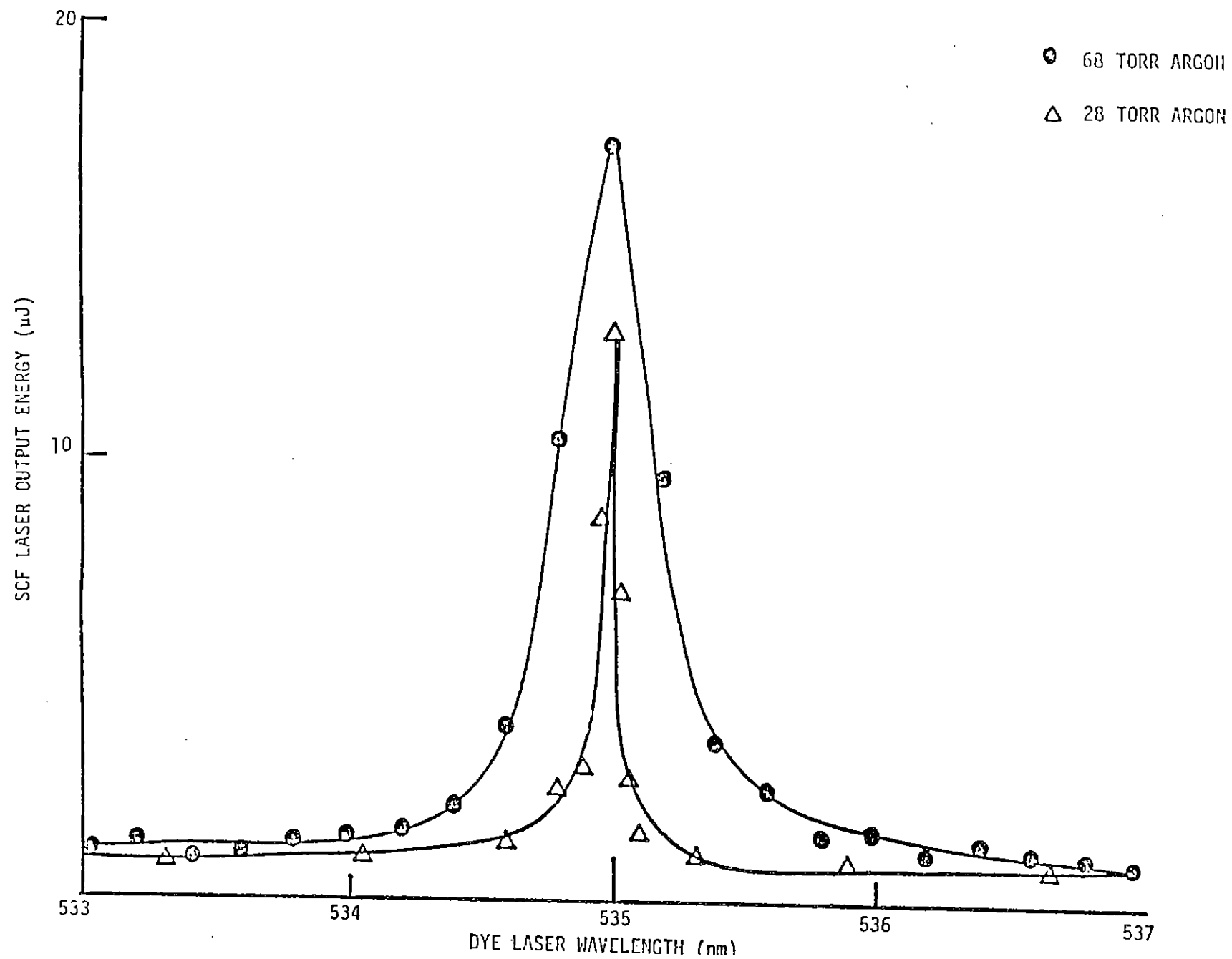


Fig. B3.13 SCF output as a function of dye laser wavelength for two buffer gas pressures with constant TICl vapour density.

shown in Figure B3.14. The latter tuning curve shows a striking asymmetry and a dip at 535 nm. Although this behaviour has been observed previously for ASR, SRS and for stimulated stokes hyper-Raman scattering, no clear theoretical model for the magnitude of the asymmetry or the dip has been proposed in the literature. The gain cross-section for near resonance Raman scattering (σ_R) is inversely proportional to the square of the detuning from resonance (B3.10-B3.13) (i.e. $\sigma_R \propto 1/\delta^2$), whereas the gain cross-section for SCF (σ_{SCF}) is proportional to $\gamma(\delta)$, the rate of inelastic collisions which populate the $7s^2S_{1/2}$ level, and inversely proportional to δ^2 . (i.e. $\sigma_{SCF} \propto \gamma(\delta)/\delta^2$). From Figure B3.13 σ_{SCF} would seem to be symmetrical about 535 nm which implies that $\gamma(\delta)$ is also symmetrical. In other words our results would indicate that while competition and possibly, to a lesser extent, interference between the SCF and ASR processes can account for the dip centred on 535 nm in the ASR tuning curve, the source of the asymmetry remains unknown. We propose to investigate the asymmetry, in more detail, in future experiments.

With 28 Torr of argon in the cell the SCF and/or ASR outputs could still be obtained when the pump laser was delayed up to 20 nsec after the dissociating ArF laser pulse. However with 68 Torr of argon, the dye laser could only be delayed by ~ 5 nsec if either an SCF or ASRL output was to be observed. This would indicate that the $Tl(6p^2P_{3/2})$ storage level population is depleted in collisions with the buffer gas atoms.

The temporal profiles of the SCF and ASR outputs were also investigated. It was found that the anti-stokes Raman laser had a pulse duration ~ 5 nsec and had a structure reflecting the peak of the dye laser pulse. The SCF output was quite different, it consisted of a single spike of ~ 1.5 nsec duration. The short pulse length of the SCF output is consistent with pulse lengths obtained in experiments on the 377.6 nm TI photodissociation laser (B3.16).

It is clear that the upper resonant level of an anti-stokes Raman up-converter will be populated, in a collisional environment, by inelastic

collisions. Although this process will obviously reduce the conversion efficiency of the anti-stokes Raman up-converter near resonance the resulting stimulated collision induced fluorescence may provide the basis for efficient up-converter; even in systems considered inefficient as Raman up-converters. The dip at 535 nm in the ASR output can be explained, in terms of competition and interference between the ASR and SCF processes or possibly Tl^*-Ar quasi-molecule formation. The source of the asymmetry in the ASR tuning curve requires further experimental and theoretical investigations.

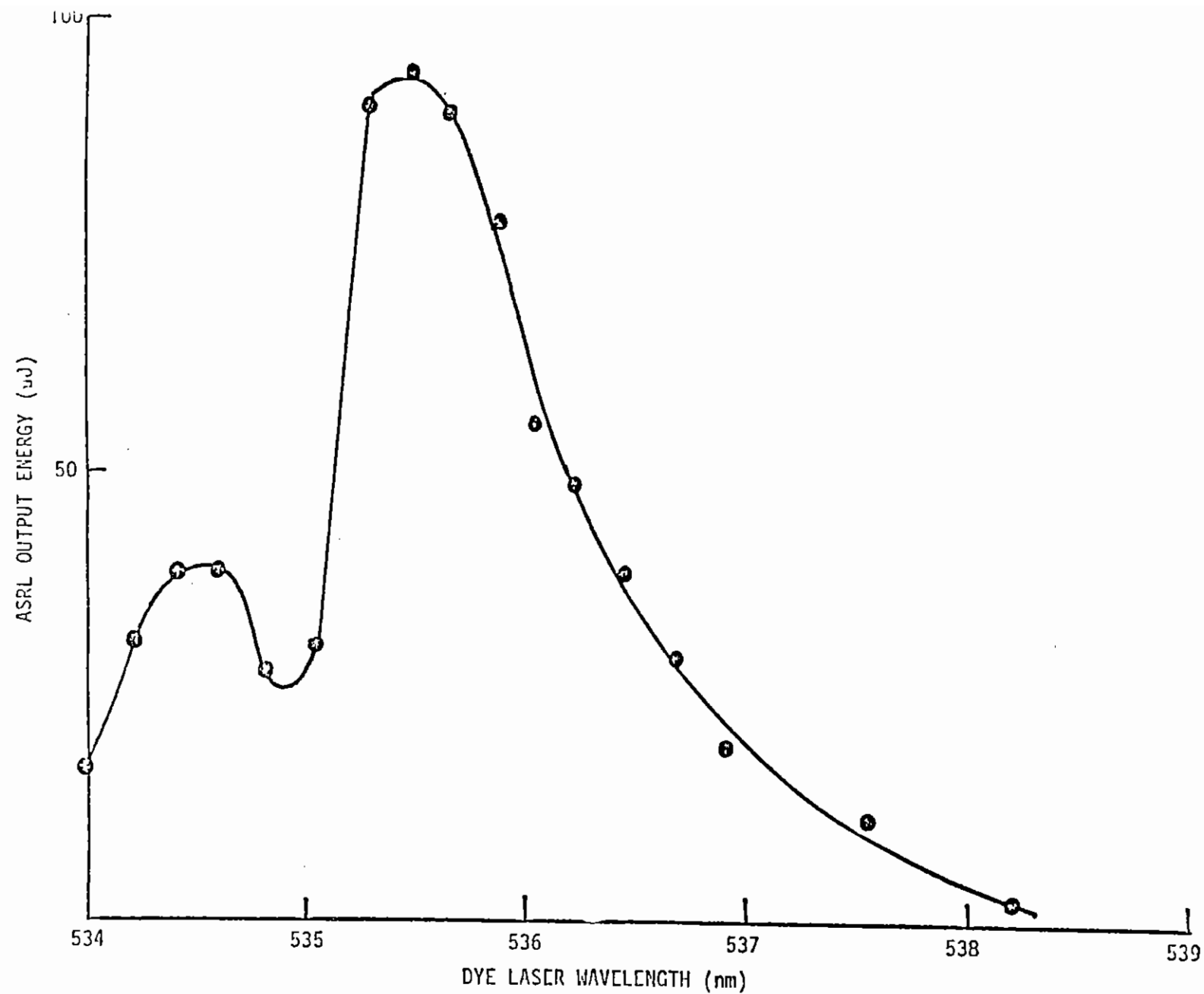


Fig. B3.14 Tuning curve of anti-stokes Raman laser₃ output as a function of dye laser wavelength with a TlCl vapour density of $\approx 3.6 \times 10^{16} \text{ cm}^{-3}$.

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B4 ENGINEERING

B4.1 Photoablation and Direct Etching of Resist Materials

G Davis, C E Webb (Oxford) and M C Gower (RAL)

Several papers have been written on laser photoablation of polymers at 193nm (B4.1-B4.6). It has been suggested (B4.3) that this phenomenon occurs as a result of the high efficiency with which polymeric bonds are broken following the strong absorption of the 6.4eV photons. The resultant energetic fragment molecules then rapidly ablate from the irradiated sites. Little experimental work has been reported on the wavelength dependence of photoablation. However, the model outlined above suggests that the efficiency of this process, at a particular wavelength, would be influenced not only by the absorption coefficient at that wavelength, but also by the probability that the absorbed photon caused bond breakage. By observation of the etching characteristics of PMMA as a function of wavelength we now present more quantitative evidence which supports this rapid bond-breaking mechanism for ablation. These results are particularly timely in view of recent observations at 308nm of the alternative thermally - induced etching process of various polymers (B4.7).

The polymer used was an experimental electron beam resist (B4.8) poly-(methylmethacrylate - co-methacrylic acid) and poly-(methyl methacrylate-co-methacryloyl chloride) (PMMA). 0.5 - 1µm thicknesses of resist were spun onto substrates of borosilicate glass. For some of the experiments the glass was coated with 1µm of chromium. This additional layer, however, did not noticeably affect the photoablative process. Lambda Physik Models 101E and 201 lasers were used for the work at 249, 222 and 193nm with pulse lengths of 20ns. The 157nm F₂ Oxford Lasers KX1 laser source produced 40mJ in an 10ns pulse. The masks used for this study were either free-standing gold structures or chromium-on-quartz masks clamped in contact with the resist. A 50cm focal length lens was used to focus the energy onto the resist and an absorption cell enabled the incident energy density to be varied. In the case of the 157nm source the energy density

was varied by altering the operating conditions of the laser. For these experiments, because of the absorption of atmospheric oxygen at 157nm, the path between the laser and substrate was evacuated to 10^{-2} torr. Experiments at longer wavelengths were performed in air.

The transmission of light through two samples of PMMA-on-quartz was measured at 249, 222 and 193nm using a D₂ arc lamp source and a spectrometer/photomultiplier combination. Assuming that the Fresnel reflection of the film is approximately that of quartz at the same wavelength, an absorption coefficient was calculated from these transmission measurements. For comparison, similar measurements were made using thin films of several other polymers that have been used in previous ablation studies (B4.1-B4.6). The absorption coefficient at 157nm was measured in a similar manner using the laser source and a pyroelectric energy meter. The measured absorption coefficients, are shown in Table B4.1. In particular, we note that of all the resists measured, at 193nm α was the lowest for PMMA.

Wavelength /nm	Absorption Coefficient/ μm					
	PMMA	PERSPEX	PET	POLYSTYRENE	PARYLENE	NITROCELLULOSE
157	5		-	-	-	-
193	0.33	0.46(0.42)	1.03	2.05	3.24	12.5
222	0.33	0.29	1.27	2.45	3.95	-
249	0.02	0.02	1.53	0.46	1.24	0.17

Table B4.1. Absorption Coefficients in μm^{-1} of six polymer used in ablation studies. The estimated errors are $\pm 0.03\mu\text{m}^{-1}$.

Results

(a) Irradiation at 193 nm

The resist was irradiated through a mask at single pulse energy ranging from ~ 1 to $150\text{mJ}/\text{cm}^2$. The exposed regions were examined in a scanning electron microscope (SEM). A plot of the fraction of resist remaining versus the single pulse laser fluence at 193nm is shown in Graph A, Figure B4.1. It can be seen that when the exposure dose exceeded $80\text{mJ}/\text{cm}^2$ complete removal of the $1\mu\text{m}$ layer of resist occurred. At lower levels of irradiation down to the threshold level of $\sim 30\text{mJ}/\text{cm}^2$, only partial ablation of the $0.6\mu\text{m}$ film of resist occurred. A typical partially etched region is shown in Figure B4.2(a). The smoothly undulating region is caused by diffraction effects from the edges of the mask. The surface texture in this region was not noticeably different from that of the unirradiated polymer. These undulations are not present at sufficiently high laser fluences such that the entire film of resist can be etched from the substrate. Contrary to other reports (B4.2,B4.5) we see no evidence of debris resulting from the deposition of ablated material either in or around the exposed region. However, we did observe deposits of ablated resist material on the quartz mask-support.

(b) Irradiation at 157nm

The experiments reported above for the 193nm source were repeated at 157nm and to our knowledge are the first reported results for 157nm laser induced photoablation. The exposure curve measured for 157nm ablation is shown as graph B in Figure B4.1. In contrast to the 193nm results of Graph A, multiple (100) pulses of the 157nm laser source were used at single pulse fluences of 1 to $30\text{mJ}/\text{cm}^2$ to obtain the data shown. It is immediately apparent from Figure B4.1 that for both 157 and 193nm irradiation there is an incident fluence threshold below which photoablation is not observed. However, since others (B4.4) have shown, that the amount of material ablated by each pulse in a multishot exposure is approximately the same, with no evidence of any accumulative effects, it is more appropriate to quote the threshold in terms of a laser intensity rather than a laser energy fluence. Hence we measure intensity thresholds for photoablation of PMMA at

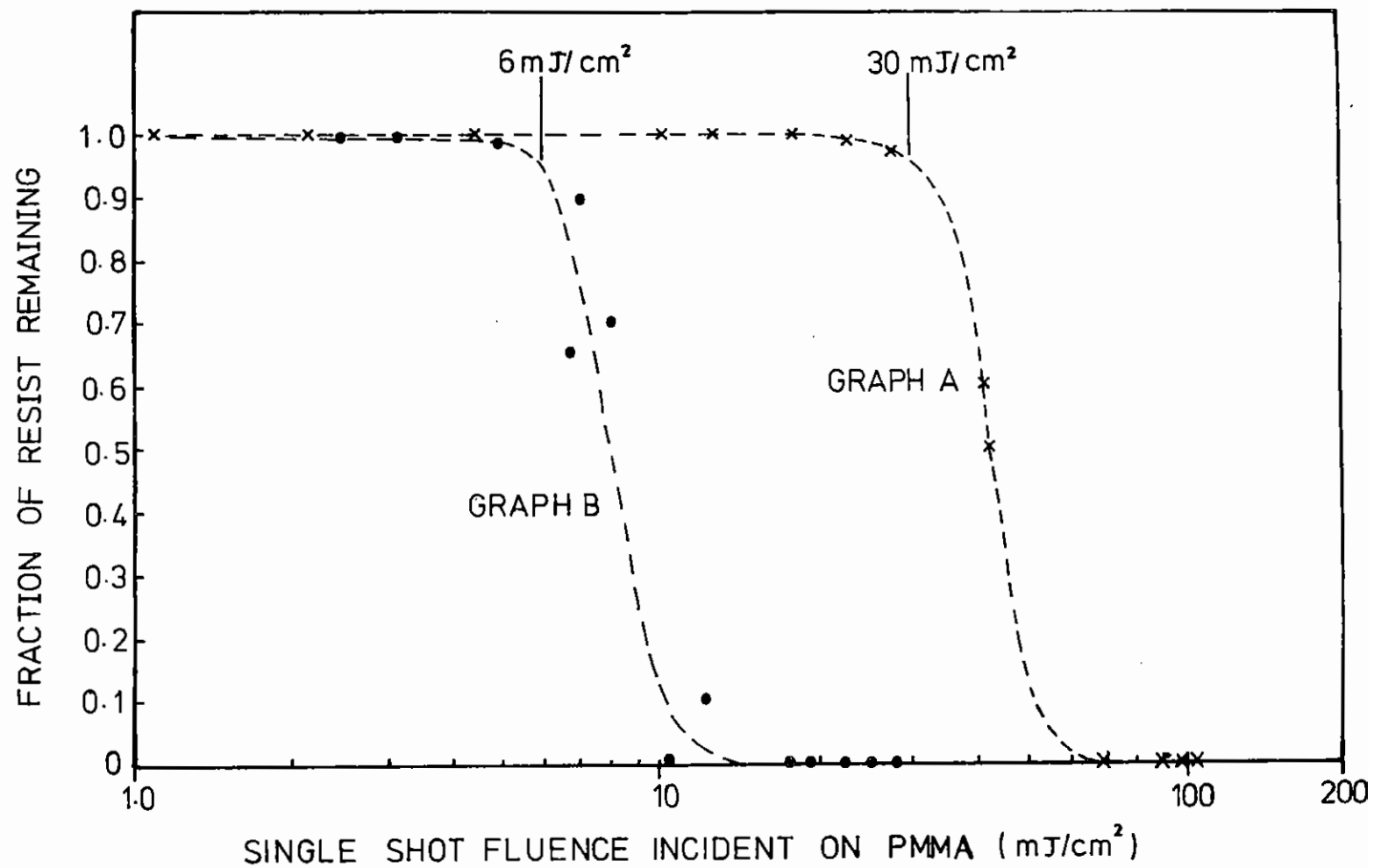
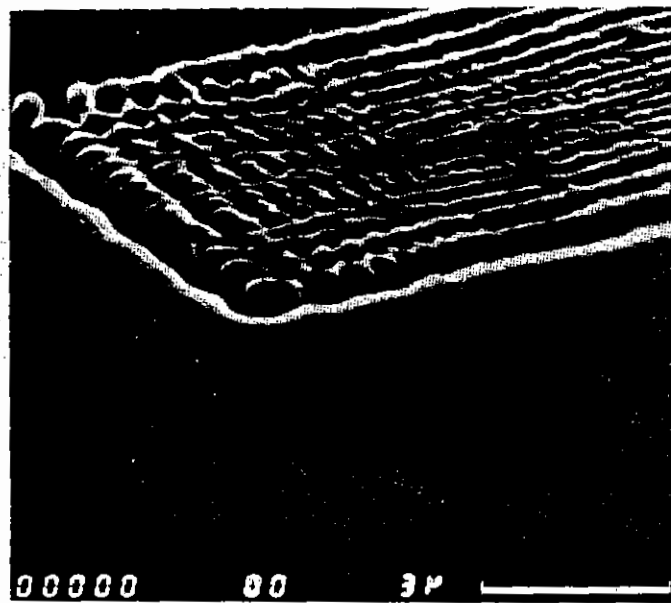
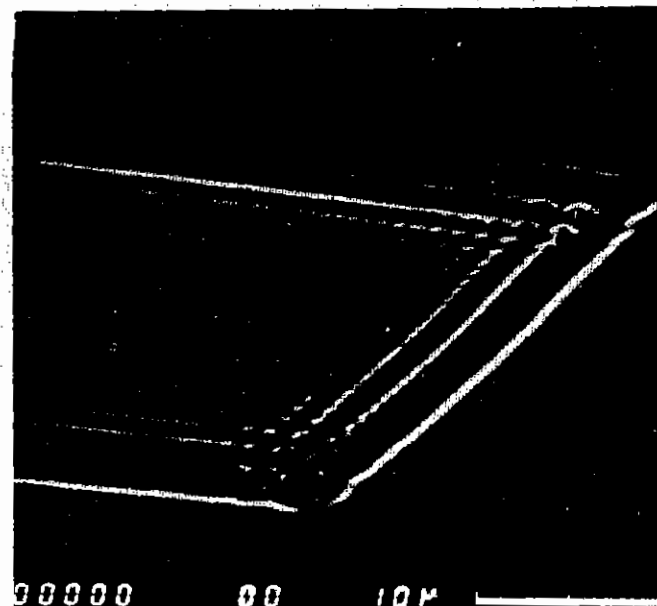


Fig. B4.1

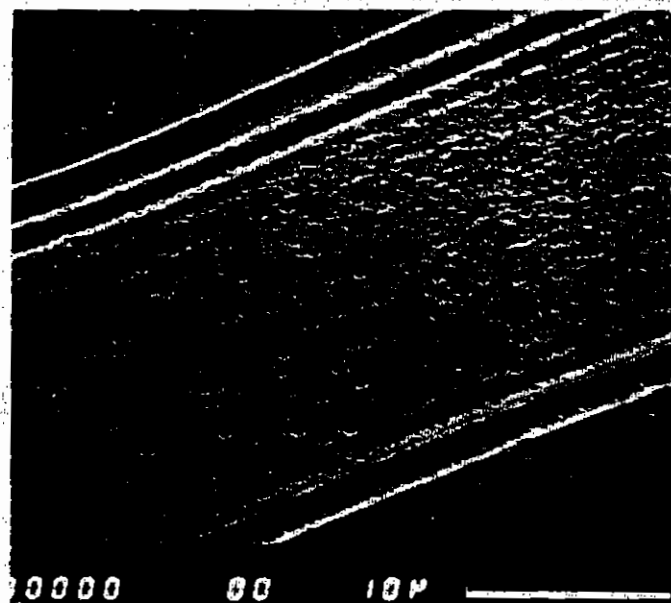
Fraction of resist remaining in the exposed region verses log (single shot fluence) for 193 (Graph A) and 157nm (Graph B) sources. The resist thicknesses were 1.0 μ m and 0.6 μ m respectively.



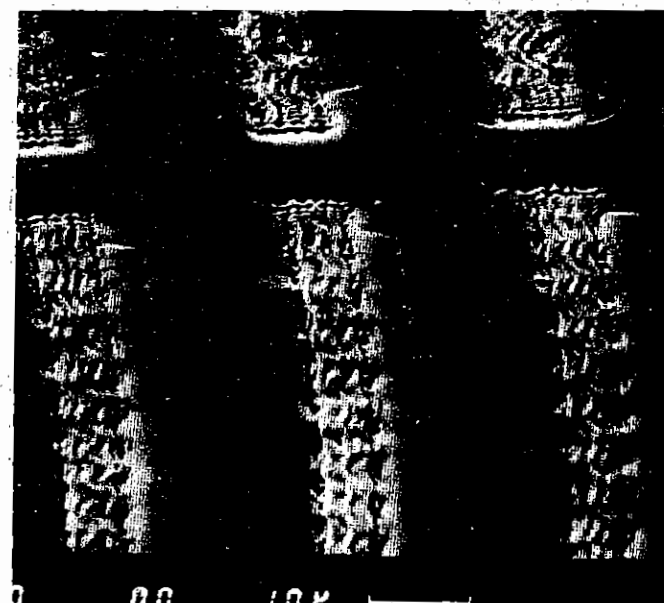
a



b

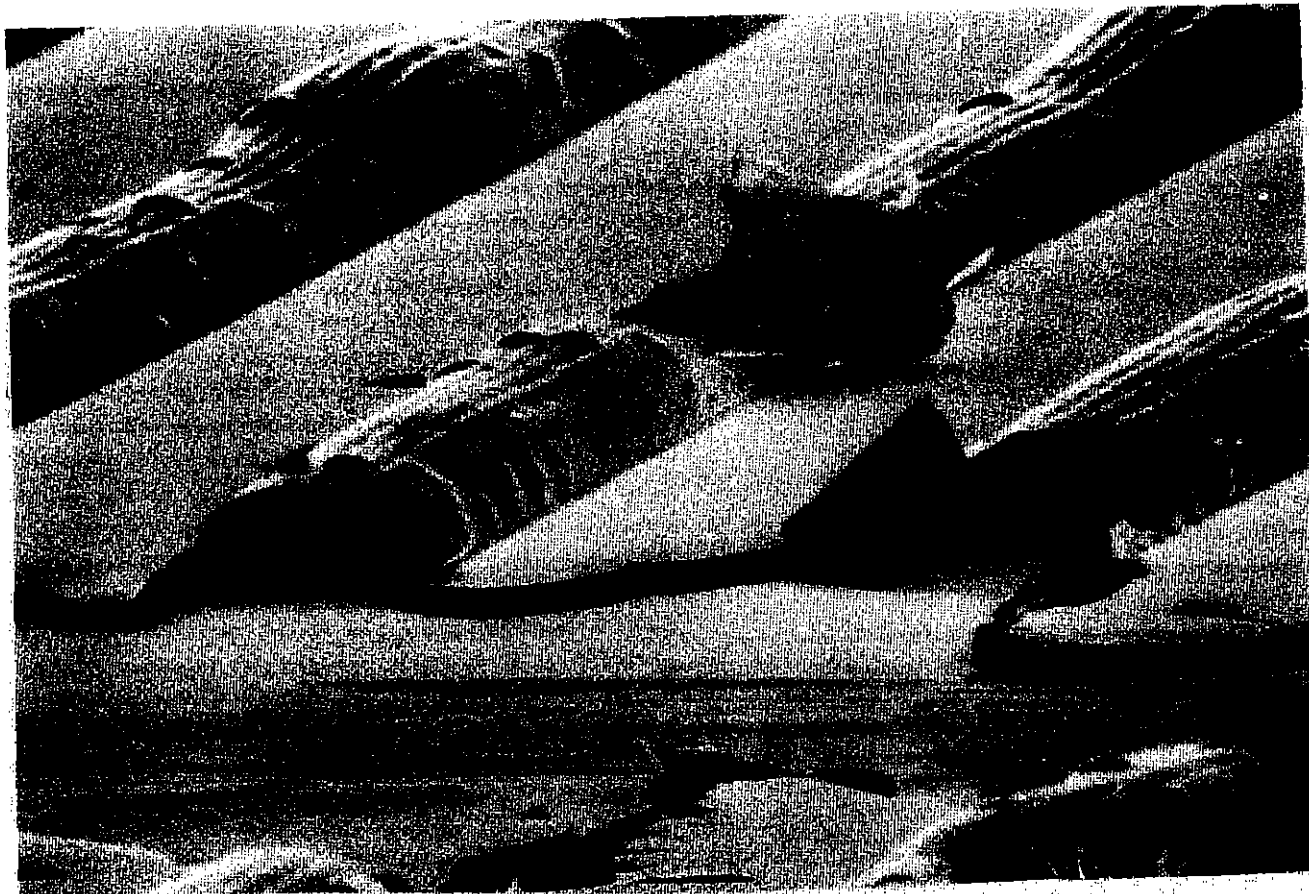


c



d

Fig. B4.2 Electron micrographs of PMMA exposures at (a) 193nm single shot at $\sim 45\text{mJ}/\text{cm}^2$, (b) 157nm 100 shots at $\sim 7\text{mJ}/\text{cm}^2$, (c) 222nm single shot at $\sim 280\text{mJ}/\text{cm}^2$, (d) 249nm single shot at $\sim 250\text{mJ}/\text{cm}^2$.



1 μ m PMMA : ArF : SINGLE SHOT

10 μ m 120 mJ/cm²

Fig. B4.3 Electron micrograph of a PMMA exposure at 193nm

193nm and 157nm of $1.5\text{MW}/\text{cm}^2$ and $0.6\text{MW}/\text{cm}^2$ respectively. The existence of this intensity threshold for observation of photoablation suggests that the rapid increase in pressure at the irradiated site due to laser-induced bond breakage must exceed a critical value.

Since the absorption coefficient of PMMA at 157nm is at least an order of magnitude larger than at 193nm, the quantum efficiency for bond scission is most likely to be larger at 157nm. Also compared to 193nm, there is an additional 1.5eV of excess energy available which can be imparted to the photofragments following bond scission by 157nm photons. Both of these effects will allow the critical pressure at which ablation begins to occur to be more readily reached at lower laser intensities. Hence, with this model of a critical pressure which must be reached for photoablation to occur, the combination of an increasing quantum efficiency for bond scission and shorter wavelength illumination are expected to lead to a lowering of the threshold intensity required for ablation. This is in agreement with our threshold intensity observations of ablation using 157 and 193nm sources. As shown in Figure B4.2(b), the quality of the photoetch obtained using the 157nm laser source was of a similar quality to the etch obtained in Figure B4.2(a) at 193nm.

From our results we estimate that the photon/monomer ratio necessary to ablate a $1\mu\text{m}$ layer of PMMA at 193nm is only ~ 0.1 . Hence ablation from the irradiated region most likely involves the primary ejection of substantial quantities of relatively large molecules in their polymeric form. Although in general samples of PMMA exposed to 193nm laser radiation responded in the manner outlined above, one sample of nominally the same resist showed a vastly different behaviour. At doses of $120\text{mJ}/\text{cm}^2$, the exposed regions had a nodular appearance and were arched upwards away from the substrate (see Figure B4.3). The reason for this inconsistency is not yet understood.

(c) Irradiation at 222nm

As is apparent from Figure B4.2(c) the effect of 222nm radiation on PMMA appears very different from that generally observed at the shorter wavelengths of 193 and 157nm. Although material has been removed from the

area irradiated at 222nm, the surface of the exposed region is pock-marked and globules of debris are visible both in and around this region. At higher energy fluences $\sim 300\text{mJ}/\text{cm}^2$, the polymer layer was almost completely removed from the glass substrate. The amount of debris was less, but the edges of the exposed regions were ill-defined and had the appearance of material that had melted. It appears that etching of PMMA at 222nm is caused by a mechanism other than direct bond scission. Since, by coincidence, the absorption depth at 222nm is similar to that at 193nm, (see Table B4.1) the enormous difference in the observed etch characteristics at these two wavelengths for comparable incident fluence is most likely to be due to a substantial reduction in the quantum efficiency for bond scission of PMMA at 222nm. Thus the critical pressure for ablation is most likely never reached. On the other hand, since the absorption coefficients are similar at the two wavelengths, a thermal model for ablation is expected (B4.7) to predict similar photoablative characteristics when similar laser energy fluences are used. This behaviour is contrary to our observations as shown in Figure B4.2(c). Most of the energy absorbed by the PMMA when exposed at 222nm is likely to degenerate into heat. The estimated rise in temperature resulting from absorption of radiation at 222nm could reasonably be expected to cause the quantity of vaporisation observed. It thus appears that for wavelengths less than 222nm, photoablation occurs by rapid direct bond scission by the laser photons which cause a correspondingly rapid rise in pressure at the irradiated site.

(d) Irradiation at 249nm

We note from Table B4.1 that the absorption depth of PMMA at 249nm is very long, $50\mu\text{m}$ and most of the incident radiation is, therefore, transmitted by the polymer and absorbed in the substrate. This resulted in the poor etching of PMMA shown in Figure B4.2(d)

In summary, we find from this work that for photoablation to be the dominant etching mechanism of PMMA it is necessary for the quantum efficiency for bond breakage to be high as well as the absorption coefficient to be $\geq 0.3\mu\text{m}^{-1}$. Our results at 157nm indicate that the absorption coefficient must not be too large since this reduces the depth of etch per shot for a

given fluence. It appears that the threshold laser intensity which is required to observe photoablation is a result of having to attain, by bond scission, a critical pressure at the ablation site. For applications such as microlithography where speed is important, a single shot process at an easily attainable fluence is more attractive. For the PMMA used in this study the optimum wavelength appears to be in the region of 193nm for which one shot at a fluence of 80mJ/cm² was sufficient to remove an 1µm layer. We have reported elsewhere (B4.9) that feature sizes of $\leq 0.8\mu\text{m}$ can be etched under these conditions.

B4.2 Spectroscopic Studies of ArF Laser Photoablation of PMMA

G M Davis (Oxford), M C Gower (RAL), T Efthimiopoulos, C Fotakis and P Argyarakis (University of Crete).

To understand further the photoablation process described above in Section B4.1 we have undertaken a spectroscopic study of the plume which is formed in front of the PMMA substrate during the ablation process. Radiation at 193nm from an excimer laser was focussed to an energy fluence of $\sim 300\text{mJ}/\text{cm}^2$ in $\sim 20\text{nsec}$ pulses onto the PMMA samples. When necessary, samples were placed in an evacuable chamber which could be pumped down to $\sim 10^2$ Torr. For time-integrated low resolution studies of the emission spectra from the plume, a Princeton Applied Research (PAR) Model 1225, 0.2m polychromator was used in conjunction with a PAR OMA2 optical multichannel analyser with a Model 1420 diode array detector. In higher resolution work a Model 1254 vidicon detector was used with a Rank Hilger monospek 1m spectrometer. For time and spatially resolved studies of the spectra at various wavelengths, different parts of the plume were sampled by moving the PMMA substrates on micrometer translators relative to the vertical entrance slit of a 0.5m-Jarrell-Ash spectrometer and an EMI 9783QR (S5 response) photomultiplier combination.

In Figure B4.4 we show survey spectra of the emission from the plume directly above the ablation site in air, vacuum and He (1 atm) environments. The most prominent bands in the region 300-600nm are due to transitions

arising from diatomic molecules. In particular, the CN B-X transition at 389nm, C₂ Swan, CH A-X transition at 431nm, CO Fourth Positive, and N₂ First and Second Positive band systems are the most readily recognisable and are characteristic features of low temperature flames (B4.10, B4.11).

It can be seen in Figure B4.4, that for many bands the emission appears stronger in the nominal vacuum environment than in air. Since the gas collision frequency is lower for the vacuum case we might expect the plume to have a much greater extent than in air at 1 atm pressure. We observed this to be the case. The emission in front of the substrate in the vacuum case was found to extend to completely fill the vacuum vessel ($\geq 25\text{cm}$ in front of the substrate). Thus the apparent increase in the intensity of the emission for the vacuum environment shown in Figure B4.4 is most likely due to an increase in the size of the emission source. Although there should be very few nitrogen atoms in our PMMA substrates, the presence of strong emission on the CN violet transition is probably due to the creation of this species by photochemically induced reactions of the ablated species with the surrounding air - even in the nominal vacuum case, the background concentration of air is sufficient to make these bands appear stronger because in this case the source was so much larger. Further evidence for this explanation of the relative intensities of the bands is also shown in Figure B4.4 by the vacuum cell was filled and evacuated with helium several times and then filled to a pressure of 1 atm with helium so that the background concentration of air should be very small. Although we observed that the visual extent of the plume was roughly the same as that observed in air, as can be seen in Figure B4.4, the intensity of all of the bands is greatly reduced. In particular, no evidence of emission from CN could be observed.

To measure the temperature of the plume immediately above the ablated site, we studied the (0,0) band of the CH(A²Δ-X²Π) transition at 431.4nm in a vacuum environment using higher resolution as shown in Figure B4.5(a). CH was chosen for detailed study since its large rotational constants make its band structure readily rotationally resolvable and its (A²Δ-X²Π) transition is spectroscopically well characterised. Furthermore, this species is one of the most likely to be produced as a direct product of the photoablation

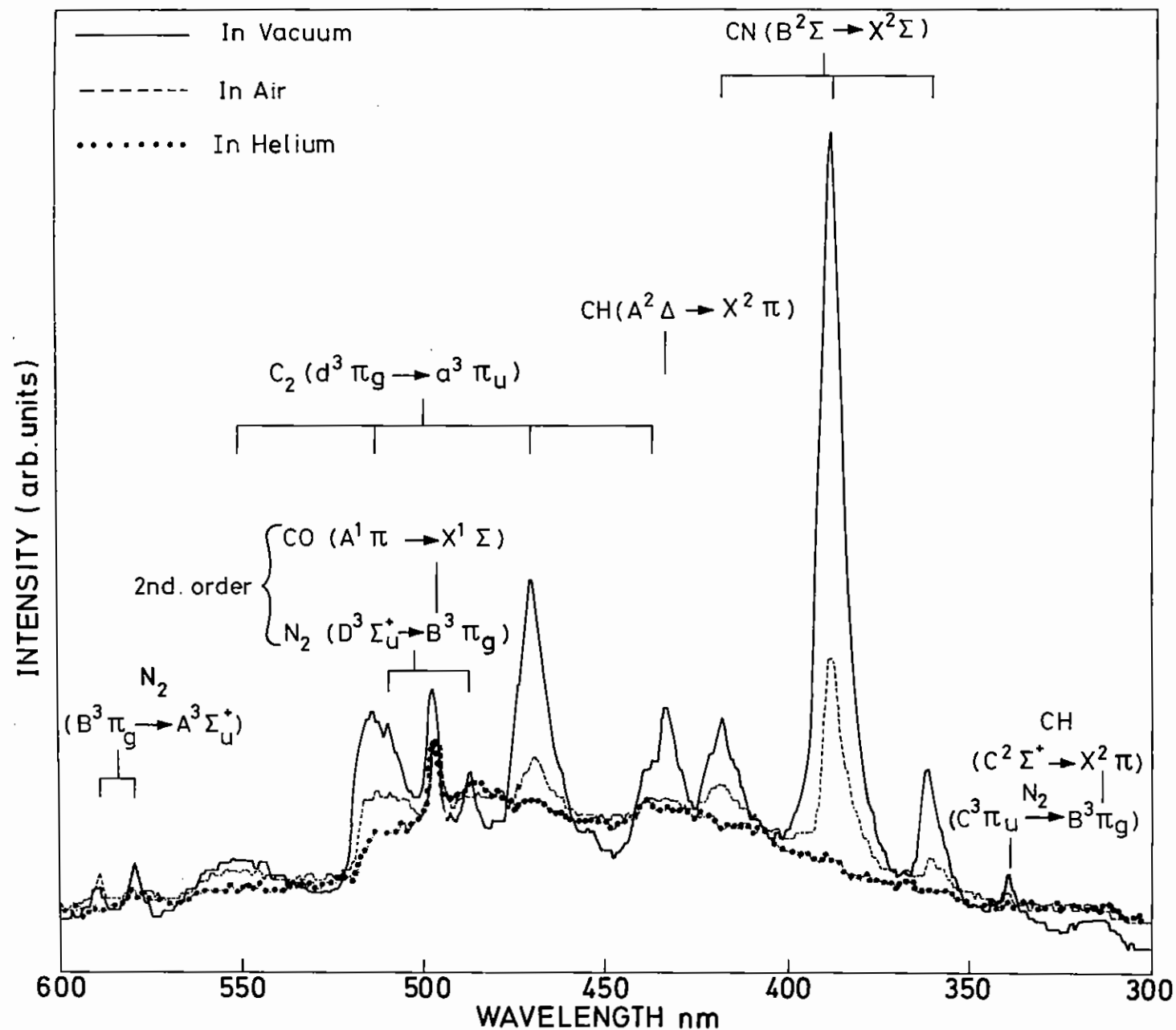


Fig. B4.4 Survey spectra of the emission between 300-600nm following irradiation of PMMA resist in air, vacuum, and He (latm) environments.

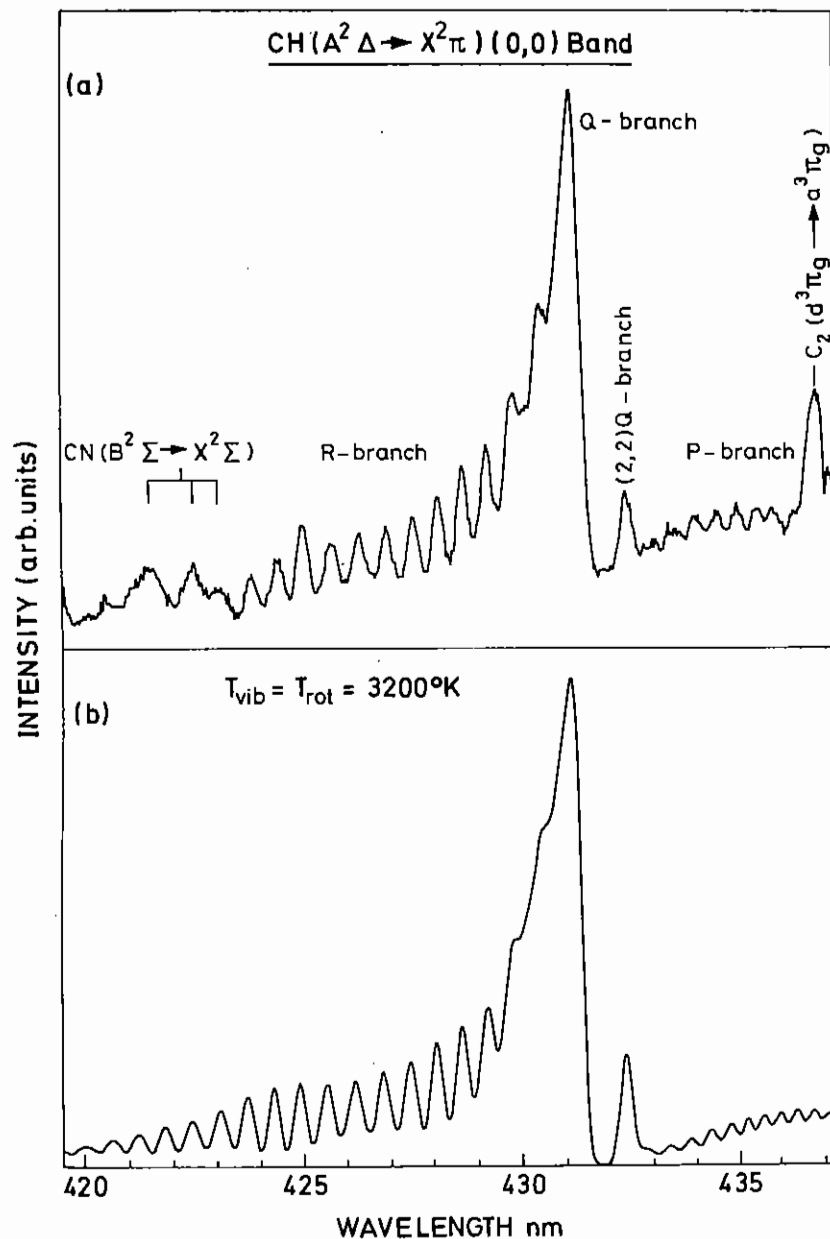


Fig. B4.5 CH(A²Δ → X²Π) (0,0) band emission.
 (a) Time integrated spectrum in vacuum environment.
 Note that bands due to CN and C₂ slightly interfere with the CH band.
 (b) Computer generated spectra with
 T_{vib} = T_{rot} = 3200 ± 200°K and a resolution of 0.3nm.

process rather than by any subsequent chemical reaction. Using appropriate line strengths for this transition, it is possible to generate a computer fit to the band by adjusting the vibrational and rotational temperatures. The best fit is shown in Figure B4.5(b) and corresponds to vibrational and rotational temperatures being in equilibrium at 3200 ± 200°K. Because of the rapid thermalisation rates between rotational and thermal motion we may then conclude that the thermal temperature above the ablation zone is also 3200°K. From measurements of the etch depth (B4.2, B4.9) and the specific heat of PMMA, we calculate that the maximum temperature obtainable when all of the laser energy is absorbed in the ablated region is only ~1900°K. It thus appears that the higher temperature which exists in the plume is a result of exothermic photochemical reactions, such as bond scission, produced by the laser photons.

To study the translational motion of some of the species present in the ablation region, we used a 15cm focal length lens to image (one to one) the horizontal plume onto the vertical entrance slits (width ~10μm) of the spectrometer and then translated the substrate. The time resolved emissions from the CH(A²Δ), C₂(d³Π_g) and CN(B²Σ) states were studied as a function of distance from the substrate. Oscilloscope traces of the CH(A²Δ) 431nm emission recorded at increasing distances in front of the substrate are shown in Figure B4.6(a)-(d). In Figure B4.6 we also show measurements taken from these traces of the time to the peak of the CH emission as a function of distance from the substrate. Provided that the quenching rates for CH are insensitive to position, then the slope of this curve provides an estimate of the translational velocity v, with which CH is moving after leaving the irradiated site. At the surface of the PMMA, we find from Figure B4.6, that the velocity away from the substrate, v₀ ≈ 4.2 × 10⁵ cm/sec for CH(A²Δ). If the mean speed of the flow v₀ = √(8RT/πM) then this speed corresponds to a translational temperature of ~11,000°K. In supersonic expansions we expect the translational temperature to be higher than the thermal temperature (B4.13) in agreement with our observations. By monitoring the time resolved emissions from the C₂(d³Π_g) and CN(B²Σ) states, in a similar manner we have measured initial translational velocities for these fragments of ~2.0 × 10⁵cm/sec and 5.1 × 10⁵cm/sec

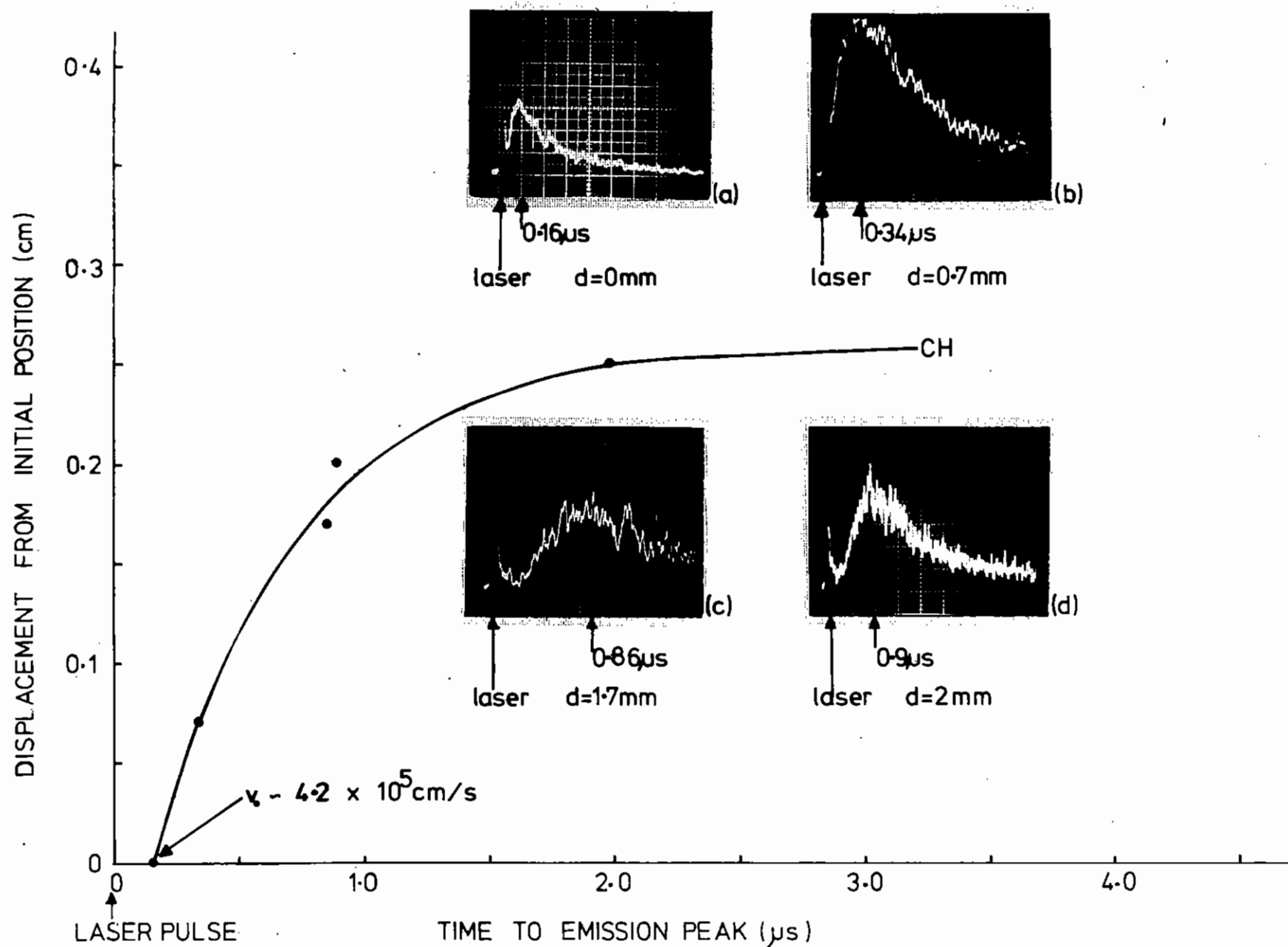


Fig. B4.6 Displacement from the PMMA substrate vs the time delay between the laser pulse and the peak in the emission from the CH(A Δ →X Σ) band at 431nm. Spectra recorded in an air environment. Oscilloscope traces of the CH(A Δ) emission as a function of position in front of the substrate are shown in Fig. 3(a)-(d). The horizontal scales were 0.2 μ sec/div for (a)-(c) and 0.5 μ sec/div for (d).

respectively.

The relatively high thermal and translational temperatures measured in this experiment show that exothermic direct bond scission reactions by the 193nm photons are responsible for the dynamics of the small fragments which occur in the plume above the ablation site. However, from other work described in Section B4.1, we know that the laser photon/monomer ratio which is necessary to ablate a $\sim 1\mu\text{m}$ layer of PMMA at 193nm is only ~ 0.1 . This implies that initially in the ablation zone relatively large molecules in their polymeric form are produced. These molecules do not appear to contribute to the emission spectra observed in Figure B4.4. It is possible that the photodissociation processes responsible for creating the emission in the plume are separate and subsequent to the breaking of the polymeric bonds which cause ablation, although it is much more likely that direct bond scission by the UV photons is also responsible for the breaking of the relatively weak polymeric bonds. The rapid rupturing of these bonds causes a rapid increase in pressure at the ablation site which then leads to the highly energetic fragments which we observe in the plume.

B.4.3 Photo-ablative etching for target fabrication

C Brown, M Davis, P Rumsby (RAL)

Ultra violet laser photoablative etching techniques have been developed considerably during the year in order to fabricate and modulate polymer targets for the glass laser scientific programme. An Oxford Lasers KX-1 ArF excimer laser has been purchased for this work which has concentrated mostly on contact printing techniques onto polystyrene shells and microdot targets. Interference techniques have also been developed in order to generate sinusoidally modulated line patterns on targets using a line narrowed KrF excimer laser. Further details of all these projects are presented in section A6.2.3.

B.4.4 Excimer laser exposure of photoresists

F Goodall, R Moody and R A Lawes (EBL Group, Technology Division, RAL)

A set of preliminary experiments was carried out using the ArF 193 nm laser on various resists which gave some interesting results and highlighted some problems specific to using a high power laser as a light source. Both ablation and normal development were investigated.

(a) Ablation

Using cross linked PMMA (XXL15) on a chrome on glass plate, ablation was observed at 50 mJ/cm^2 , however with the same resist on a silicon wafer and a patterned aluminium maske (from Plessey, Towcester) no ablation effects were observed at this power density.

Using AX1350 J Shipley positive photoresist no ablative effects were observed until about 100 mJ cm^2 .

(b) Normal Development

Using A21350 J, 193 nm at 50 mJ/cm^2 the resist appeared to give a negative resist behaviour, the pattern staying on where it was illuminated and with a slightly stronger developer the unexposed resist was removed. This is a very interesting and potentially useful result. This behaviour is of course known when exposing to e-beams, but the exposure dose required is high which tends to preclude its use on economic grounds.

Using A22350 J and A22400, exposures were made at lower doses at 193 nm and 308 nm. The main problem is to try and achieve an exposure in a single shot. Clearly a dose greater than the ablation dose cannot be used so the total integrated dose required must be less than this. Only AX2400 at 308 nm seemed to be near this goal, developing out in two shots and nearly out in one shot at 50 mJ/cm^2 in normal 3.5:1 developer. It is hoped that a slight increase in dose and developer concentration would reach the target at 308 nm. Resist work clearly needs to be done to achieve both single shot exposures at lower wavelengths.

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B5 UVRF Facility Development

M C Gower (RAL)

In response to the increasing demand for time on the UVRF and the somewhat overcrowded working conditions, in September 1983 the facility obtained extra full time support of an additional person (Mr G Hogg), and enlarged to include an additional laboratory which is adjacent to the original UVRF laboratory. At present, this new area houses the excimer laser (Lamda Physik 101E) pumped dye laser (FL 2002E) system. Thus experiments which only require the use of the dye laser or the line-narrowed excimer laser can be scheduled for simultaneous operation. In this way the facility has been able to substantially increase the number of experiments scheduled. The two laboratories can also operate in tandem for those experiments which require the use of both laser systems.

In order to obtain high output energies and more reliable operation, the Lambda-Physik FL2001E dye laser was upgraded to a FL2002E system in September 1983.

A high energy F₂ laser (Oxford Lasers (KX-1) which produces ~40mJ of energy in a ~10nsec pulse in the vacuum ultraviolet spectral region at 157nm was added to the facility in December 1983, and its use is shared with the Target Fabrication Group at RAL. This laser can also operate on the rare gas fluoride laser wavelengths and produces ~ 350mJ at 193nm at 1pps. for example. To overcome the timing jitter problems encountered in the past, the Lamda Physik 101E line-narrowed oscillator/EMG 201 amplifier system has been upgraded to a Model 150ETS line-narrowed injection locked system which typically produces ~500mJ in a 25nsec pulse at 249nm with a line width of ~0.3cm⁻¹ at a maximum repetition frequency of 10pps. While the single switch of this device eliminates timing jitter problems it also has the additional advantage that it can be used for line-narrowed operation at 193 and 351nm.

An extra Princeton Applied Research OMA2 console has been purchased by the facility so that there are now independantly operating OMA systems available

for use (a vidicon and a diode array system), in both laboratories.

While the facility was fully operational for greater than 90% of the scheduled time, the few operational problems encountered throughout the year were mainly associated with the thyratron switches in the lasers. It is hoped that these problems can be minimised in the future by stocking in the laboratory a comprehensive set of spares for the lasers.

C HIGH POWER KrF LASER DEVELOPMENT

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Editor F O'Neill

SECTION C

INTRODUCTION

F O'Neill (RAL)

The long term aim of the high power KrF laser development work is to obtain 100 J, 1 ns UV laser pulses from the "Sprite" system. This will be achieved using two stages of optical multiplexing viz a KrF multiplexer followed by a methane Raman multiplexer. Design work on a KrF x8 multiplexer which will deliver 200 J in 8 ns is well advanced and all the optical components and beam lines have been carefully itemised and positioned on plan drawings. Ordering of equipment is now progressing at a steady rate the plan being to actually run the multiplexer in about year.

In addition to the above design work we are also carrying on some more fundamental studies particularly in the areas of UV optical damage and KrF-pumped Raman generators and amplifiers. This year the work on Raman amplifiers has been particularly productive and we have carried out our first large scale experiments using the Sprite laser and a segmented mirror to pump a methane wave-guide. The plans for this experiment were described in detail in last years report. The results have been very encouraging. In the Sprite-pumped CH_4 wave-guide we have succeeded in obtaining a small signal gain of e^7 at the first Stokes wavelength (268 nm) and have amplified a probe beam from 1 mJ up to 2 J representing a power extraction efficiency of ~ 15%.

While we are very pleased with this level of success in these preliminary Raman experiments we have been hampered by the low energy (~ 1 mJ) and large line-width (0.3 cm^{-1}) of the 268 nm first Stokes probe beam. The low energy means we do not saturate the amplifier and the large line-width keeps the gain low. As described in section C2 below we have rectified this situation and the Raman generator can now produce pulses of > 10 mJ in a line-width ~ 0.15 cm^{-1} . New experiments with the large wave-guide are now being planned in which we hope to demonstrate > 50% extraction efficiency and amplified first Stokes beams of diffraction limited beam quality. These are

of course the requirements which must be met if KrF-pumped Raman amplifiers are to be used successfully in a laser pulse compression scheme.

C1 KrF MULTIPLEXER DESIGN

D Craddock, S Hancock, E Madraszek, F O'Neill, I N Ross and M J Shaw (RAL)

In the 1983 Annual Report a conceptual design was shown of a 5 x 12 staged multiplexer for compressing the 60 ns Sprite pulse to a Raman shifted pulse of 1 ns duration and > 100 J energy. Since that time measurements of the Sprite e-beam pump pulse have shown it to have a duration of approximately 65 ns and in order to maximise the multiplexer efficiency it has been decided to adopt an 8 x 8 multiplexing array giving an overall compression of x 64 and an output pulse duration still ~ 1 ns.

In this past year design work has progressed on the x 8 KrF multiplexer. Detailed design of the x 8 Raman multiplexer must await the outcome of the wave guide experiments which are still on-going and these are described in section C5 below. Present plans are to install the KrF multiplexer on Sprite during the 1985 reporting year to provide a 200 J, 8 ns output for target work. The design of the KrF multiplexer can be understood by referring to the block diagram shown in Figure C1.1 and the laboratory diagram shown in Figure C1.2. The multiplexer is driven by a 100 mJ, 8 ns KrF oscillator pulse obtained using an electro-optic switch fitted to a Lambda Physik EMG150 laser. This beam is relayed through an overhead helium-filled beam pipe into the multiplexer room. In the relay pipe the beam is spatially filtered and is magnified up by a factor of two to have dimensions 20 mm x 40 mm.

On entering the multiplexer room the beam is periscoped down to working height and is cut into two beams each of dimensions 20 mm x 20 mm and energy ~ 50 mJ. One of these beams is delayed by 48 ns and then both beams are split into six using dielectric coated beam splitters. The output from this system consists of twelve, 8 ns pulses with a delay of 8 ns between each pulse. These pulses are then amplified sequentially in the e-beam-pumped KrF pre-amplifier GOBLIN over a time period of 96 ns. After amplification

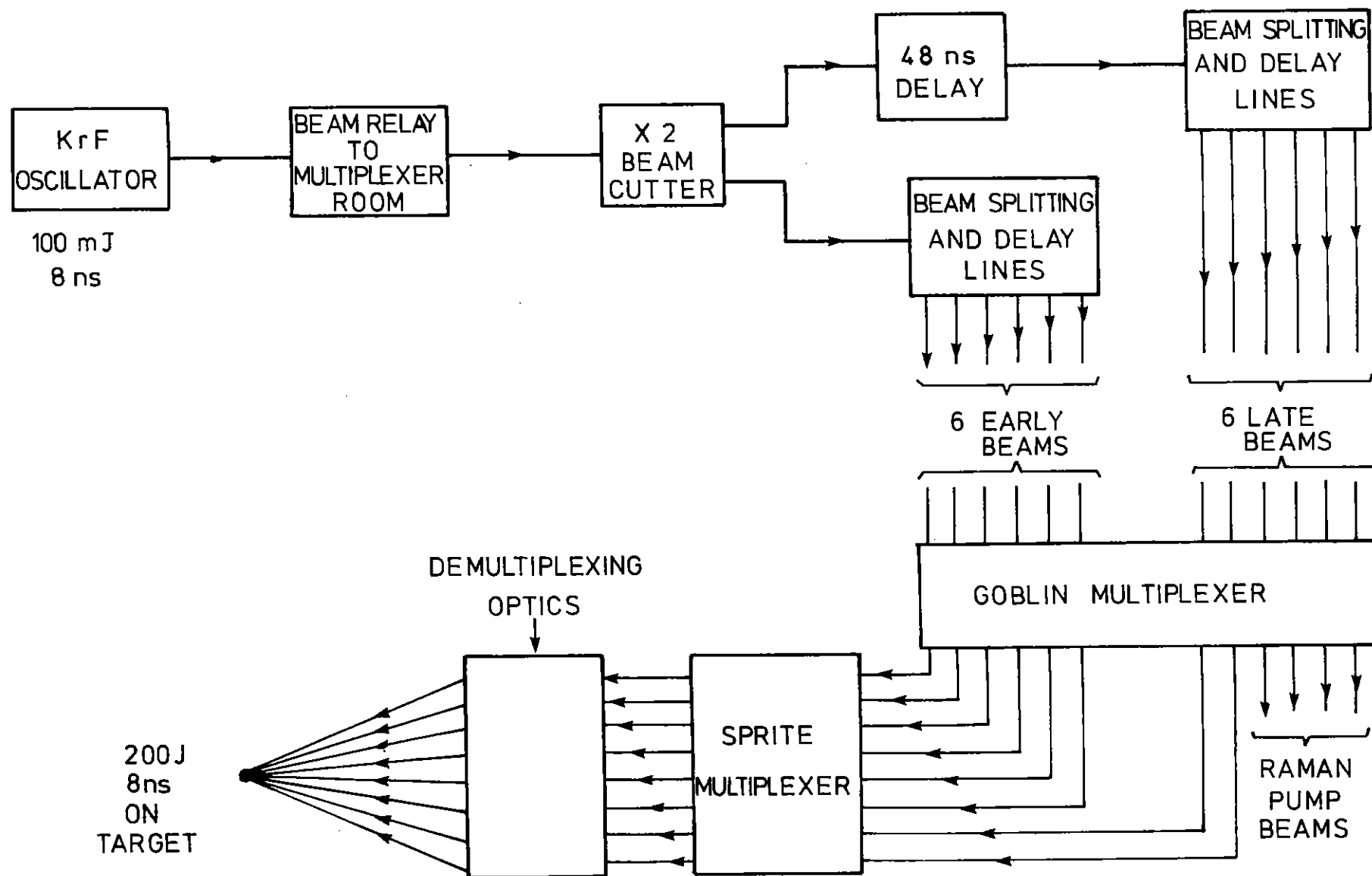


Fig. C1.1 Block diagram of x8 KrF multiplexer.

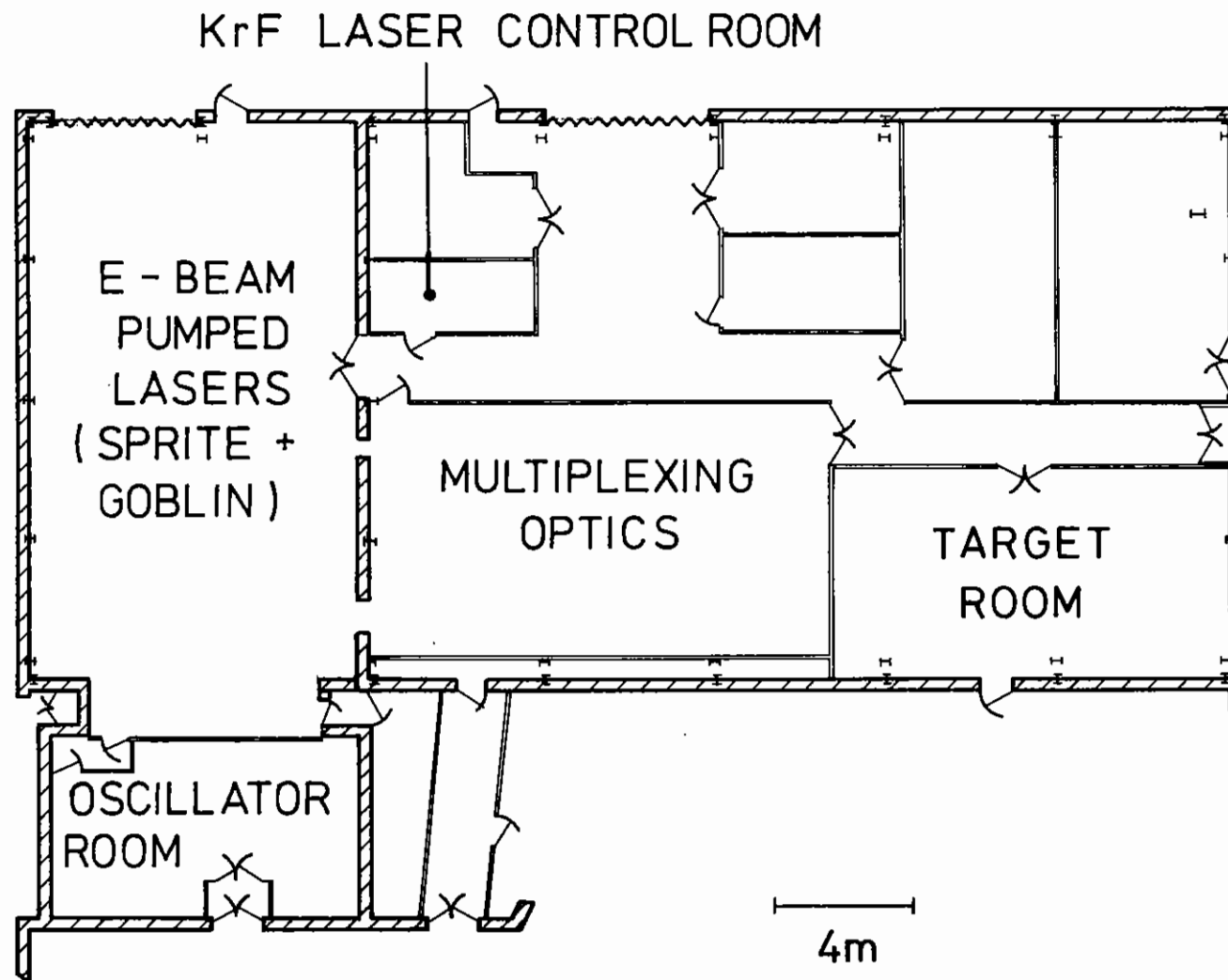


Fig. C1.2 KrF multiplexer laboratory area.

the energy in each beam is ≈ 1 J and the beam size is still ≈ 20 mm.

After amplification in GOBLIN four beams are split off and will later be used to pump Raman generators. The remaining eight beams are then amplified sequentially in Sprite to provide eight output beams, each with approximately 25 J of energy. The Sprite multiplexing optics also increase the beam size from 20 mm to 80 mm to avoid optics damage problems. Finally the optical delays between the pulses are removed by de-multiplexing optics to give a laser energy of approximately 200 J in 8 ns for target experiments.

The actual path taken by one of the KrF beams is shown in detail in Figure C1.3. The total distance travelled by the beam from the oscillator room to the target room is ≈ 100 m and in an effort to minimise distortion the beams are enclosed in pipes over much of their path. The beam pipe from the oscillator room to the multiplexer room is helium-filled and contains optics to magnify the beam and carry out spatial filtering. The designs of the beam pipes on Goblin and Sprite are shown in Figures C1.4(a) and (b). In both cases the systems have been designed to minimise the angles between beams thus maximising the geometrical extraction efficiency of the amplifiers. Also of-course the amplifiers are double passed in telescope arrangements so that costs can be kept down by using the smallest possible multiplexing optics while still being able to extract the large aperture amplifiers in an efficient manner.

The input plate of the Goblin beam pipe holds 24 convex lenses (focal length 250 cm) one lens for each input beam and one for each output beam. In the Goblin beam pipe the 20 mm beams are expanded gradually over a distance of 10 m to fill the 80 mm laser amplifier aperture. A 7.5 m radius of curvature mirror at the back of Goblin recompresses each beam back down to its original 20 mm size at the output of the beam pipe. A spatial filter pinhole is used in the output leg of each beam path to provide beam clean-up as necessary.

The input plate of the Sprite beam pipe holds eight, small diverging lenses (focal length ≈ -112 cm) for the 20 mm input beams and eight large diverging

lenses (focal length ≈ -600 cm) for the 80 mm output beams. The laser fluence is maximum on the large lenses and is ≈ 500 mJ/cm². The input beams are expanded slowly over a distance of 13.8 m to fill the 260 mm Sprite aperture and a 17 m radius of curvature mirror recompresses the beams down to 80 mm diameter on the output lenses. This change of beam size is necessary to avoid damaging the output lenses. With this optical arrangement the geometrical extraction efficiency of Sprite is $> 95\%$ on the second (saturating) pass and a total output of > 200 J should be easily achieved. After amplification the 25 J KrF beams are directed to the demultiplexer array to remove delays and give a 200 J, 8 ns pulse for target work.

Work on the x 8 KrF multiplexer is progressing steadily. All the laboratory areas have now been constructed and taken over. Most of the components for the Goblin preamplifier have been delivered and construction work is underway. The optical and mechanical design of the multiplexer hardware has been completed and purchase orders have been placed for the optical components (lenses and mirrors). Orders for mechanical components (beam pipes, mirror mounts, mirror stands, etc) will be placed in the next few months. The only major component not yet designed is the spatial filter/beam pipe from the oscillator room to the multiplexer room and this design work will be completed shortly. It is expected that the multiplexer will be operational in approximately one year.

C2 KrF PULSE GENERATOR DEVELOPMENT

E M Hodgson and I N Ross (RAL)

The pulse generator section of the KrF laser system consists of a KrF oscillator and preamplifier based on a commercial Lambda Physik EMG150 laser together with a Raman generator. The broad requirements on pulse generation are:

- (1) An injection pulse of about 100 mJ in 8 ns at 249 nm for the KrF multiplexed laser. This pulse should have a divergence of less than 100 μ rad. For initial multiplexing of Sprite, line narrowed operation is

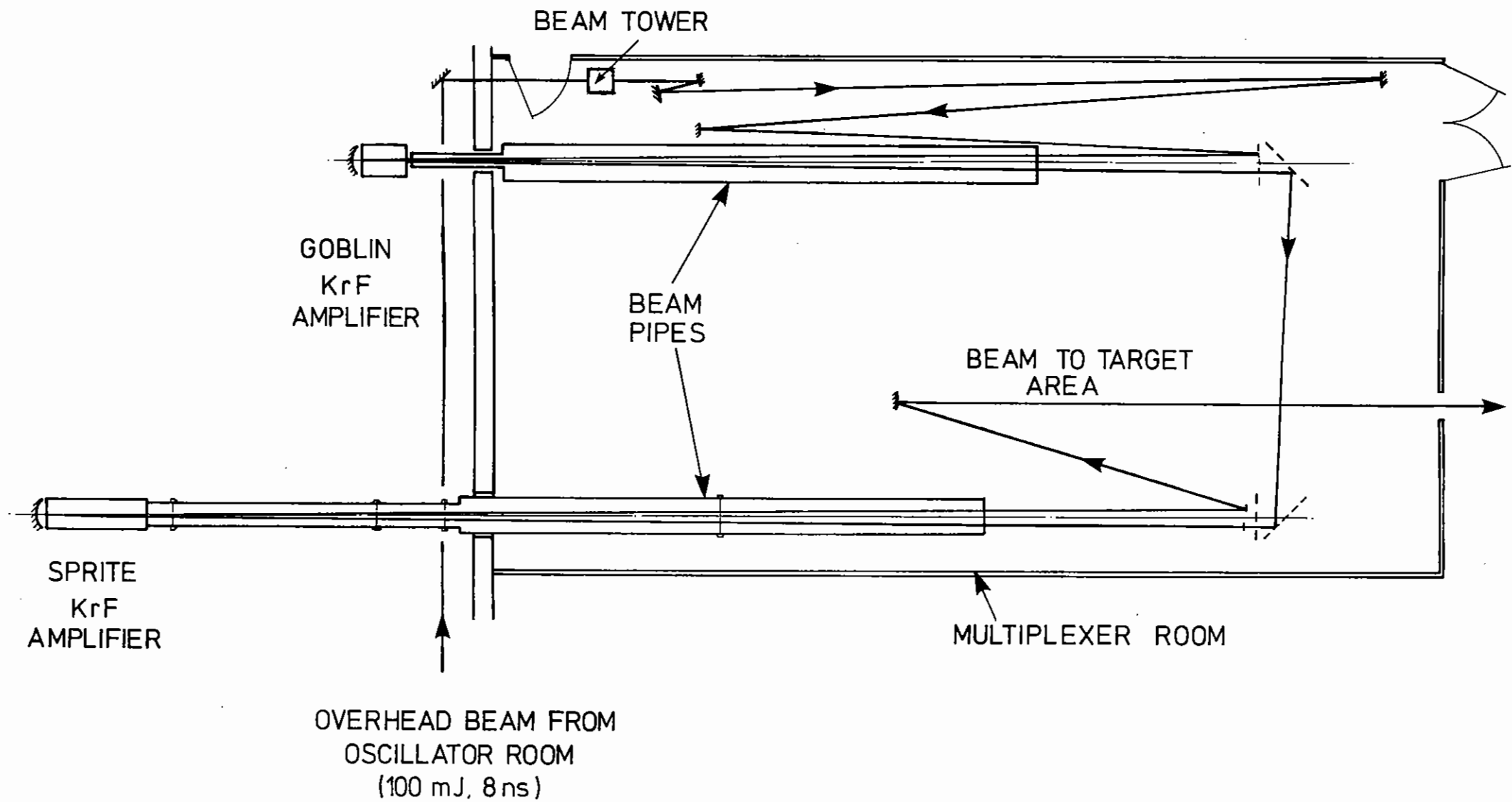
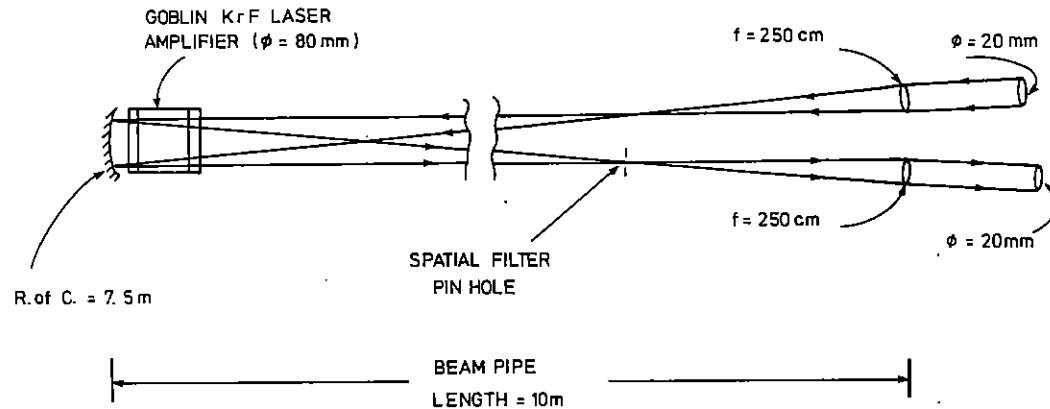
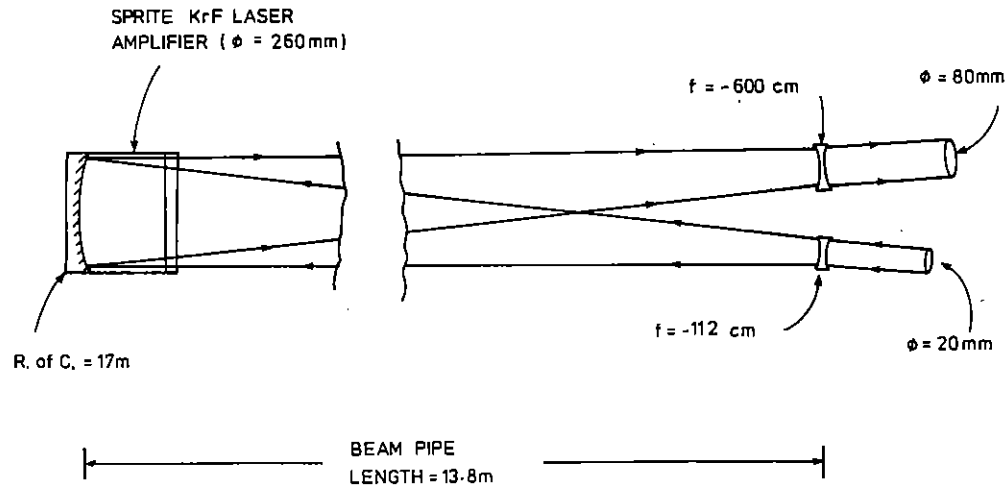


Fig. C1.3 KrF laser beam path through the Goblin and Sprite multiplexing optics.

(a) GOBLIN BEAM PIPE



(b) SPRITE BEAM PIPE



not required; however for the eventual staged multiplexer a line-width of $< 0.1 \text{ cm}^{-1}$ will be used.

(ii) A Raman shifted pulse of about 100 mJ at 267 nm in a linewidth less than 0.1 cm^{-1} for injection into a KrF pumped forward Raman amplifier.

The following paragraphs present some development work directed toward these goals.

C2.1 Oscillator Development

The principal limitation of the existing oscillator was its spectral linewidth which at about 0.4 cm^{-1} was too broad for maximum gain in a methane Raman generator or amplifier where the pump linewidth should be much less than the Raman linewidth of 0.3 cm^{-1} . The shot to shot stability of the oscillator linewidth was also not satisfactory. Figure C2.1(a) shows the existing resonator optics in which all the dispersion is achieved using beam expanding prisms coupled to dispersing prisms. Using this number of prisms is lossy and cannot generally achieve its design performance since it places extreme requirements on the optical quality and stability of components in the expanded beam. Further, the output beam although diffraction limited in the vertical sense is elongated in the horizontal sense by virtue of the prism dispersion, giving an output with wavelength changing along this direction. Figure C2.1(b) shows the modified resonator optics. A combination of fewer prisms and an etalon was proposed to achieve the 0.1 cm^{-1} linewidth, while a higher reflectivity output mirror helps to bias the system towards a greater number of passes through the dispersing optics. The prisms give a lower beam expansion factor of 20 and a lower loss ($\approx 50\%$ in double pass without antireflection coatings). They should give a linewidth of $\approx 0.5 \text{ cm}^{-1}$. The solid etalon with a free spectral range of 1.0 cm^{-1} and a finesse of about 10 reduces this linewidth to $< 0.1 \text{ cm}^{-1}$ without a large reduction in output power.

In practice this scheme has worked successfully. Figure C2.2 shows Fabry-Perot interferograms of the oscillator output with and without the intracavity etalon giving linewidths of 0.5 cm^{-1} and 0.15 cm^{-1} respectively

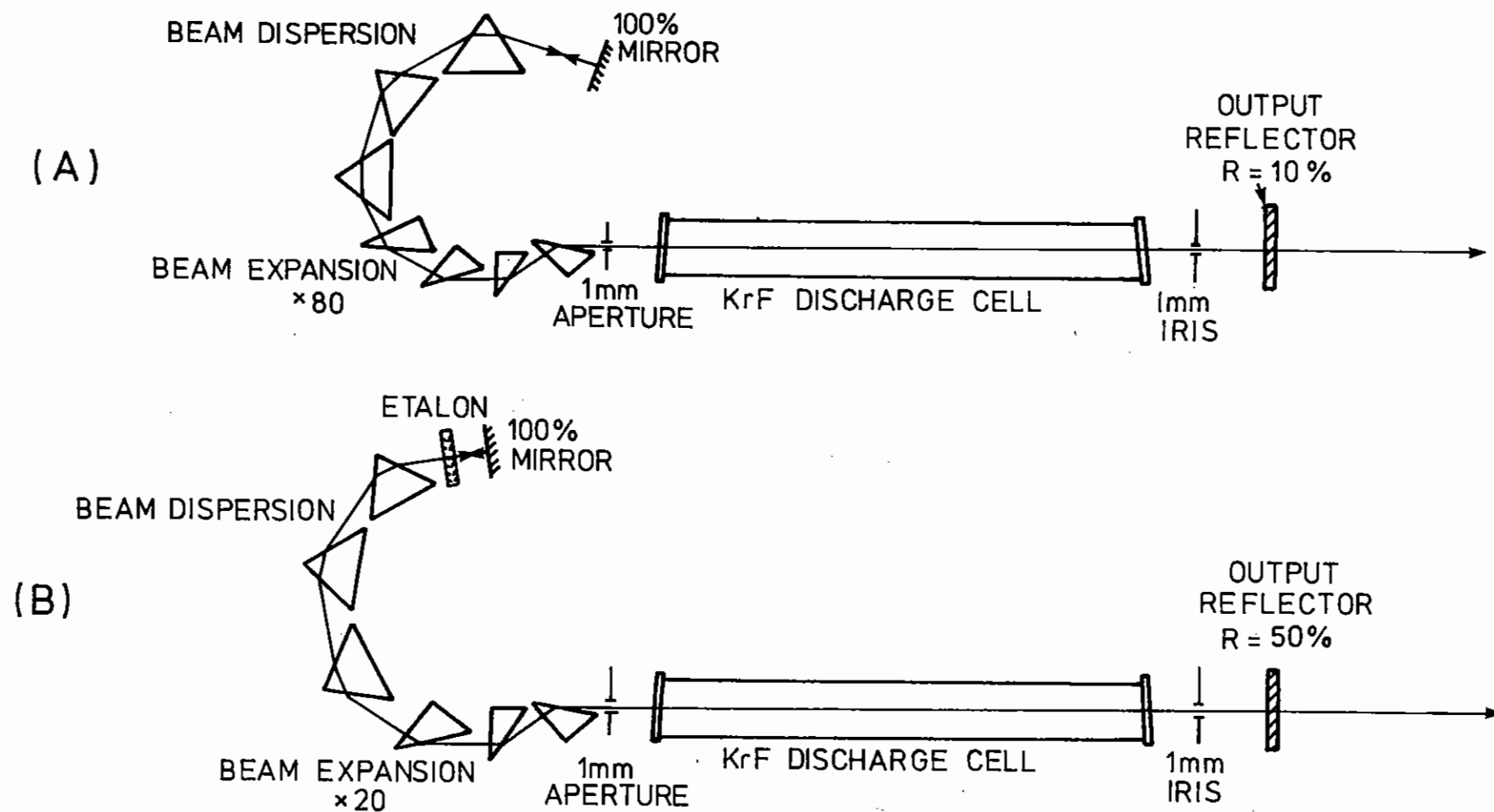
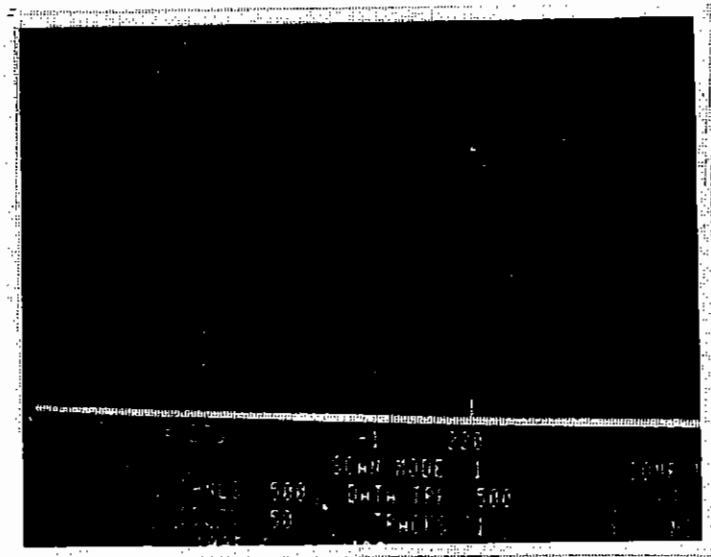


Fig. C2.1

Lay-outs of line-narrowed oscillator optics.
 (A) optics as supplied by manufacturer,
 (B) new design.

LINE - NARROWED KrF OSCILLATOR

(A) WITHOUT
ETALON



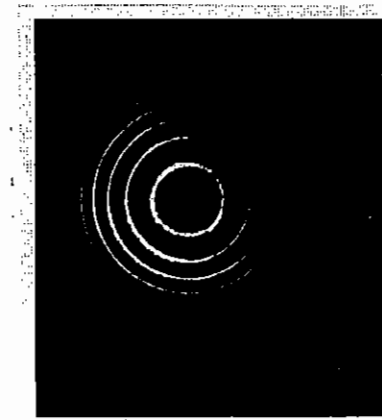
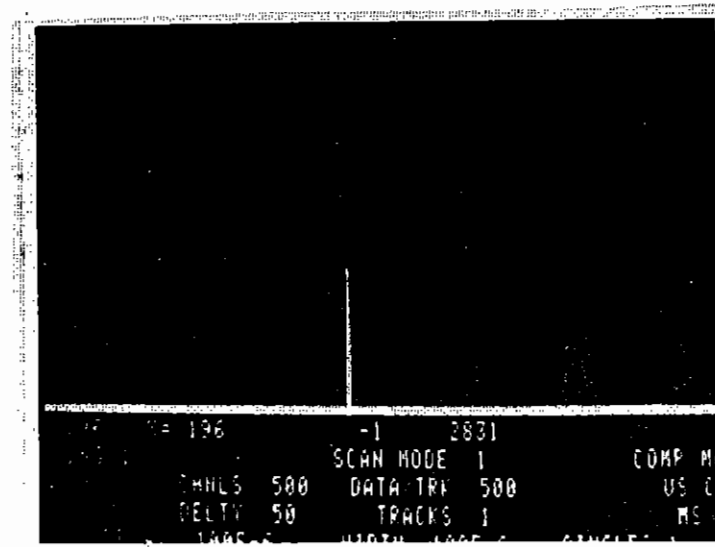
O.M.A. SCAN



FABRY - PEROT
FRINGES

$$\Delta\bar{\nu} = 0.5 \text{ cm}^{-1}$$

(B) WITH
ETALON



$$\Delta\bar{\nu} = 0.15 \text{ cm}^{-1}$$

Fig. C2.2 Fabry Perot interferograms showing output spectrum of oscillator.

before deconvolution of the instrumental linewidth. The shot to shot variation in linewidth is very small and not a problem. However good mechanical and temperature ($< 0.5^{\circ}\text{C}$ with the present arrangement) stability are required to prevent both a change of output beam alignment and the appearance of satellite spectral components separated from the main component by the free spectral range of the etalon.

The beam distribution with etalon inserted is now largely diffraction limited in both vertical and horizontal senses showing a circular distribution. The output energy with the etalon is $\approx 100 \mu\text{J}$ (@ 30 kV) which is the same as was obtained with the original cavity. However the beam intensity in the far field is considerably higher by virtue of the reduced horizontal beam divergence.

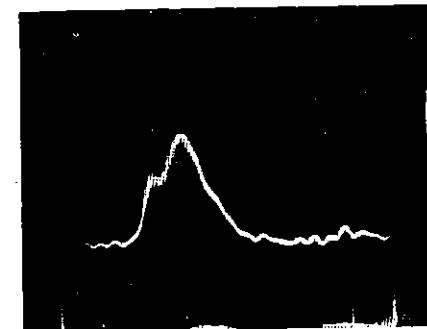
Some observations on the beam temporal profile are noteworthy. There is broadstructure (Figure C2.3) with peaks corresponding to two, four and six passes of the resonator. The output mirror is adjusted to maximise the second peak (it does not affect the first peak at all) which has considerably better spectral and spatial characteristics as compared with the first peak. There is also significant finestructure on a subnanosecond scale showing modulation levels of some tens of percent. This we believe is due to beating between modes which are confined to a small number by the degree of line-narrowing present. It should be noted that such structure may be present with broader spectral output but would then be on a smaller time scale and below the limit of time resolution of our oscilloscope. This structure on our beam can be of considerable importance in nonlinear processes such as Raman generation; see paragraph C2.3.

C2.2 Preamplifier Studies

In the existing EMG150 system the amplifier is operated in the injected unstable oscillator mode with the injected beam undergoing around five passes before emerging as the output beam. This mode is energy efficient, provides maximum duration pulses and gives adequate locking efficiencies. Alternative modes of operation have been considered which may be less critical on timing and alignment, which may give higher power in shorter

BROAD STRUCTURE

{ SCOPE BANDWIDTH = 100 MHz
SCALE = 10 ns/cm }



FINE STRUCTURE

{ SCOPE BANDWIDTH = 1 GHz
SCALE = 5 ns/cm }

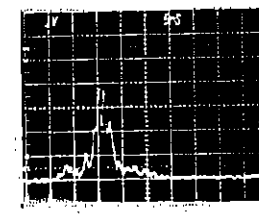


Fig. C2.3 Oscilloscopes showing the temporal structure of the oscillator output pulse.

pulses and which may lead to reduced beam divergence. In particular some experimental assessment of the double pass amplifier mode has been made.

(i) Amplified spontaneous emission (ASE) measurements were made. Spatial nonuniformities in the single pass ASE were found to originate from stray reflections from internal components in the amplifier. These together with the Fresnel reflections from the end windows form a feedback channel for parasitic oscillations. These parasitics lead to significant depletion even in single pass. Heavy saturation of ASE occurred under double pass conditions, showing that double pass amplification of even low level signals should be sufficient for efficient energy extraction.

(ii) Figure C2.4 shows the saturation curve for the amplifier in single pass. This gave a value for the saturation intensity of 1.0 MW/cm^2 and a peak single pass output energy of 100 mJ.

(iii) Figure C2.5 shows the scheme used for double passing the amplifier. The performance did not depend critically on timing. Several double pass geometries were tried with very similar results.

Energies of 250 mJ were obtained in 10 ns at a peak power of 25 MW. More than 80% of the energy was within 90° rad. All beam measurements were taken after a spatial filter to eliminate the ASE. It was found that there is a negligible amount of ASE within a 300° rad cone and it is easily eliminated using a spatial filter without attenuating the output beam. Figure C2.6 shows oscilloscope traces of the temporal profile of the pulse component and ASE component respectively under conditions of optimum amplifier injection pulse timing. An interesting possibility available with a double pass scheme is demonstrated by Figure C2.7 which shows the output pulse evolution for different optical delays between the amplifier and retro-mirror. The saturating return beam quenches the gain in the amplifier and so prematurely cuts off the amplification of the forward going beam. The result is an output pulse of reduced pulse duration where this pulse duration is twice the optical delay between amplifier and retro-mirror.

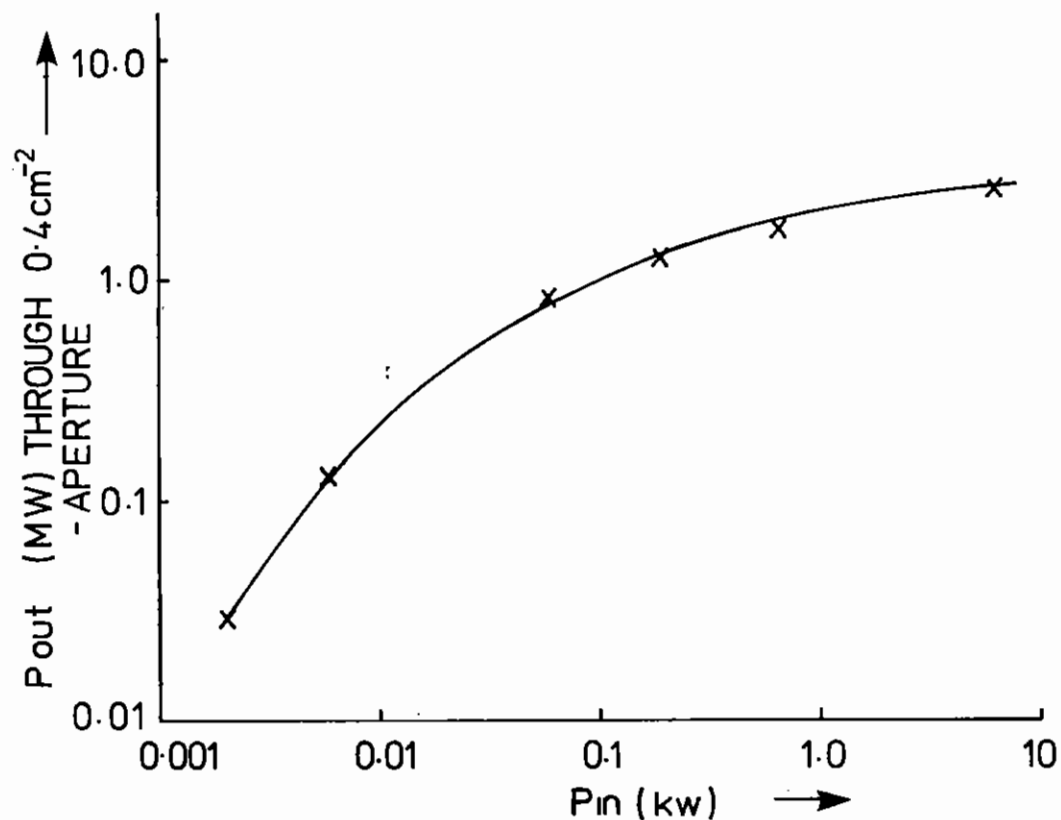


Fig. C2.4 Saturation curve for the KrF gain medium with a line-narrowed input.

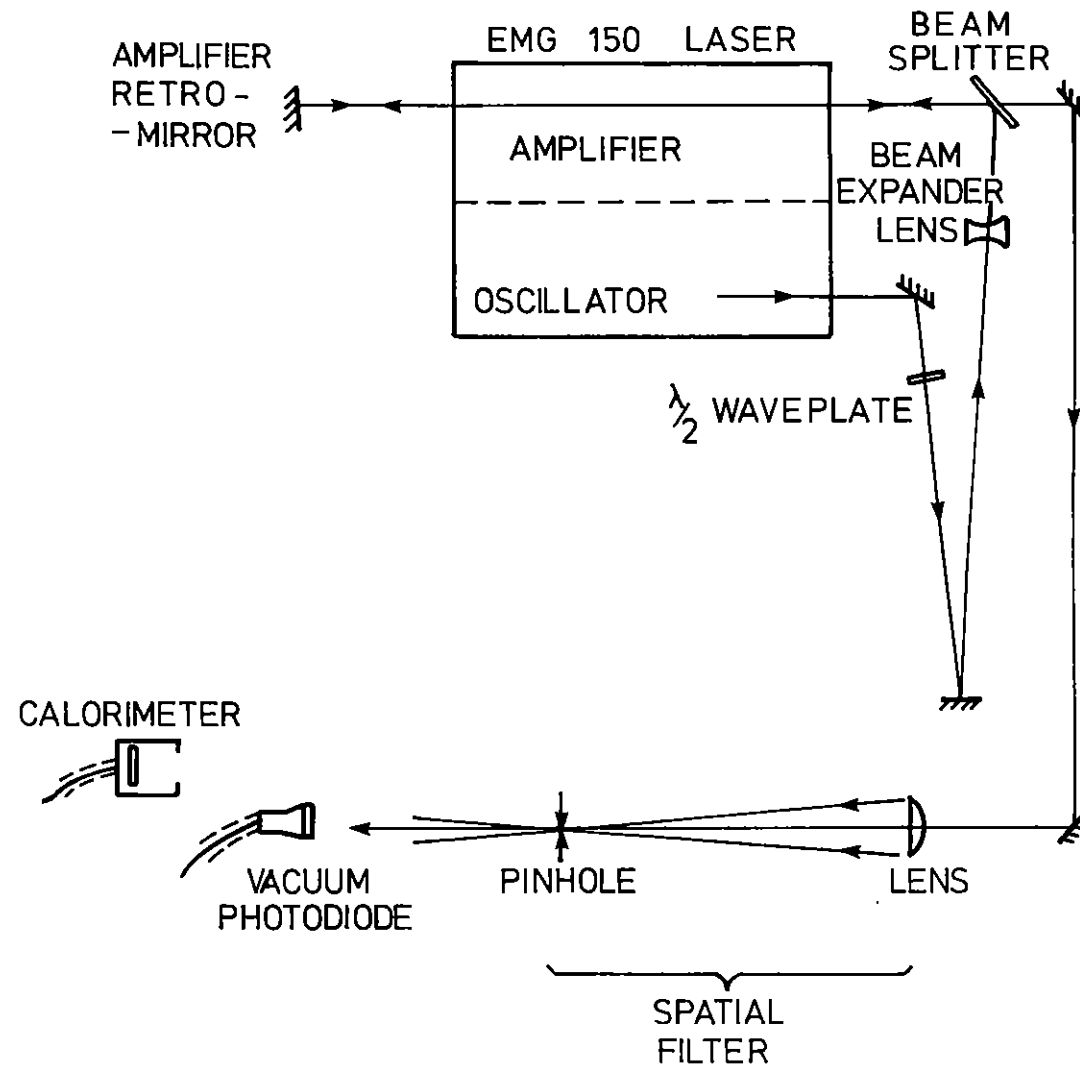


Fig. C2.5 Scheme for double passing the laser amplifier.

PULSE ONLY
(ASE BLOCKED)



HORIZONTAL SCALE = 10 ns/cm

A.S.E ONLY
(PULSE BLOCKED)

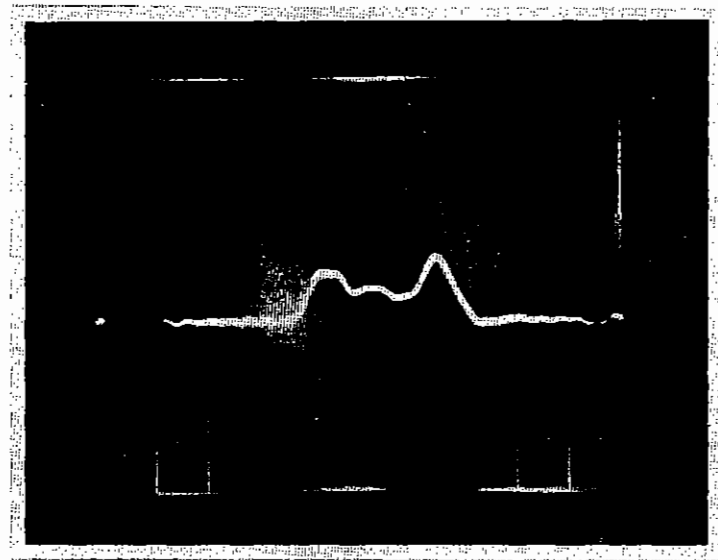
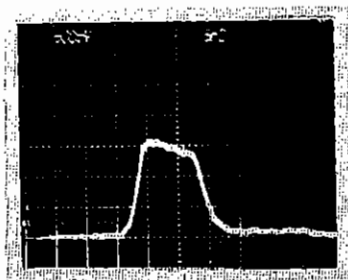


Fig. C2.6 Oscilloscope showing the temporal profile of the amplified oscillator pulse (upper) and the ASE output from the amplifier (lower).

OPTICAL DELAY
IS 4ns IN
DOUBLE - PASS



OPTICAL DELAY
IS < 0.5ns IN
DOUBLE - PASS



Fig. C2.7

Oscillograms of laser pulses obtained by amplifying the oscillator pulse on a double pass through the amplifier for two different optical delays on the second pass.

C2.3 Backward Raman Generation

To meet the future requirements for Raman shifted pulses it is necessary to achieve conversion efficiencies of the order of 50% from KrF fundamental (249 nm) to first Stokes in methane (268 nm). This appears not to be possible with a forward Raman generator due to competition with higher order Stokes and anti-Stokes conversion processes. Following the work of Tomov et al (C2.1) who achieved 80% conversion to first Stokes in the backward direction and showed that it has phase conjugate properties we have set up a backward Raman generator using methane gas, as shown in Figure C2.8. The major requirement for efficient backward Raman conversion is that the linewidth of the pump source must be as narrow as possible and certainly $\ll$ the Raman bandwidth (0.3 cm^{-1} in this case). Hence our specification of a bandwidth of 0.1 cm^{-1} .

The arrangement shown in Figure C2.8 was used to investigate the dependence of backward and forward Raman processes on the methane gas pressure and on the focusing conditions of the pump beam. Figures C2.9 plots the photon conversion efficiency into the backward first and second Stokes components as a function of gas pressure. No other Raman components or the Brillouin component were detected in the backward direction. A peak efficiency of 40% was obtained, considerably less than that achieved by Tomov and perhaps due to the somewhat broader linewidth of our pump source. At the optimum pressure the efficiency of the first backward Stokes was measured for several focal length pump beam focusing lenses. Figure C2.10 shows a peak at a focal length of 2 m (input beam size $\sim 15 \text{ mm}$) when the focal intensity was just below the breakdown threshold of the gas.

Figure C2.11 shows typical time resolved pulses of the pump, first backward Stokes for two different focal lengths and second backward Stokes (at high gas pressure). The pulse compression process inherent in the backward Raman process when operating at high efficiency is very evident especially in the second Stokes component. This second Stokes was typically just a single pulse of duration < 0.5 nanoseconds with a peak power a factor > 2 higher than the pump pulse (measurement limited by detector and scope resolution). It is noted that the pulse duration is limited to $> \approx 0.1 \text{ ns}$ by the

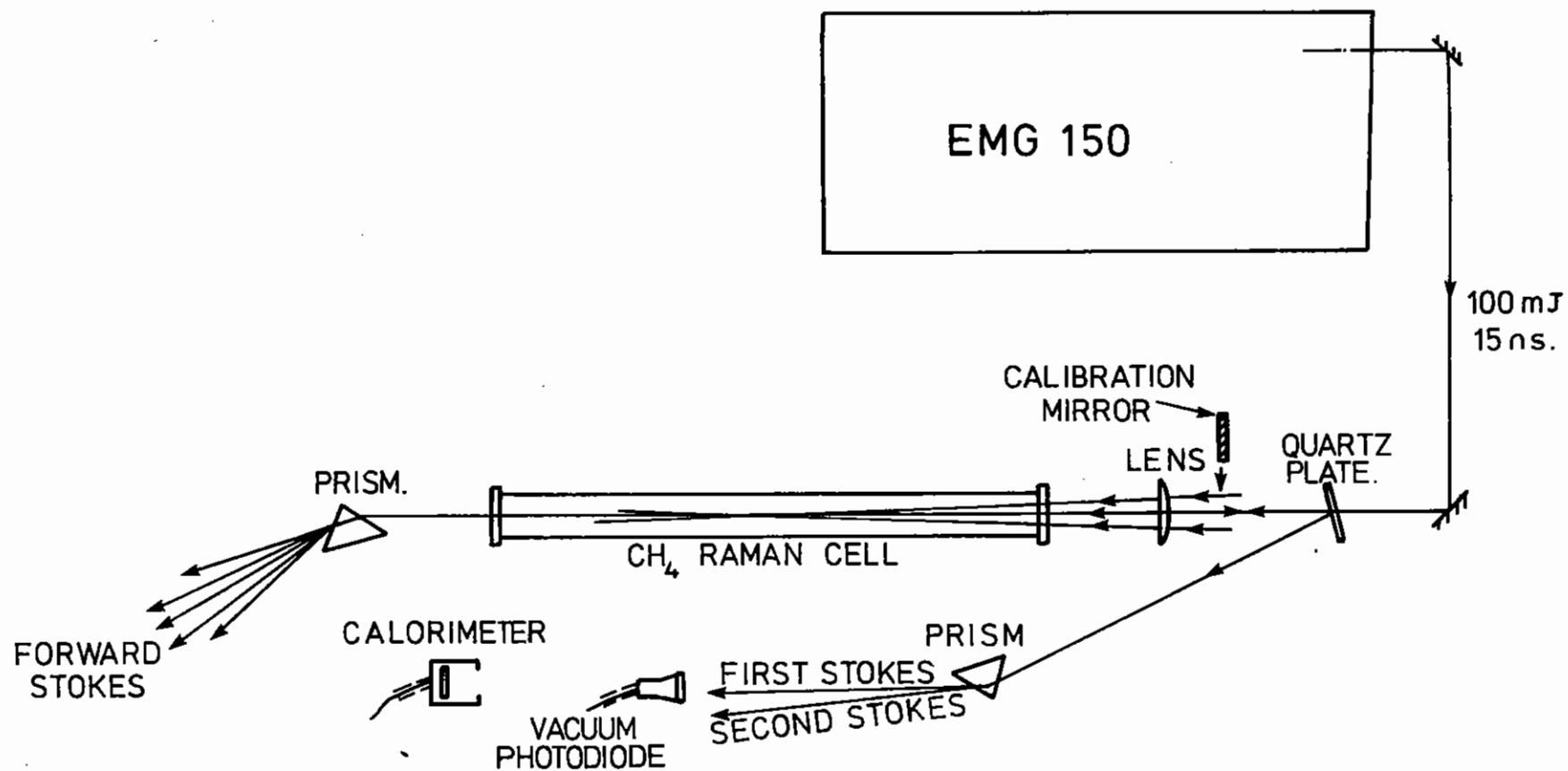


Fig. C2.8 Experimental arrangement used for the backward Raman generator experiment.

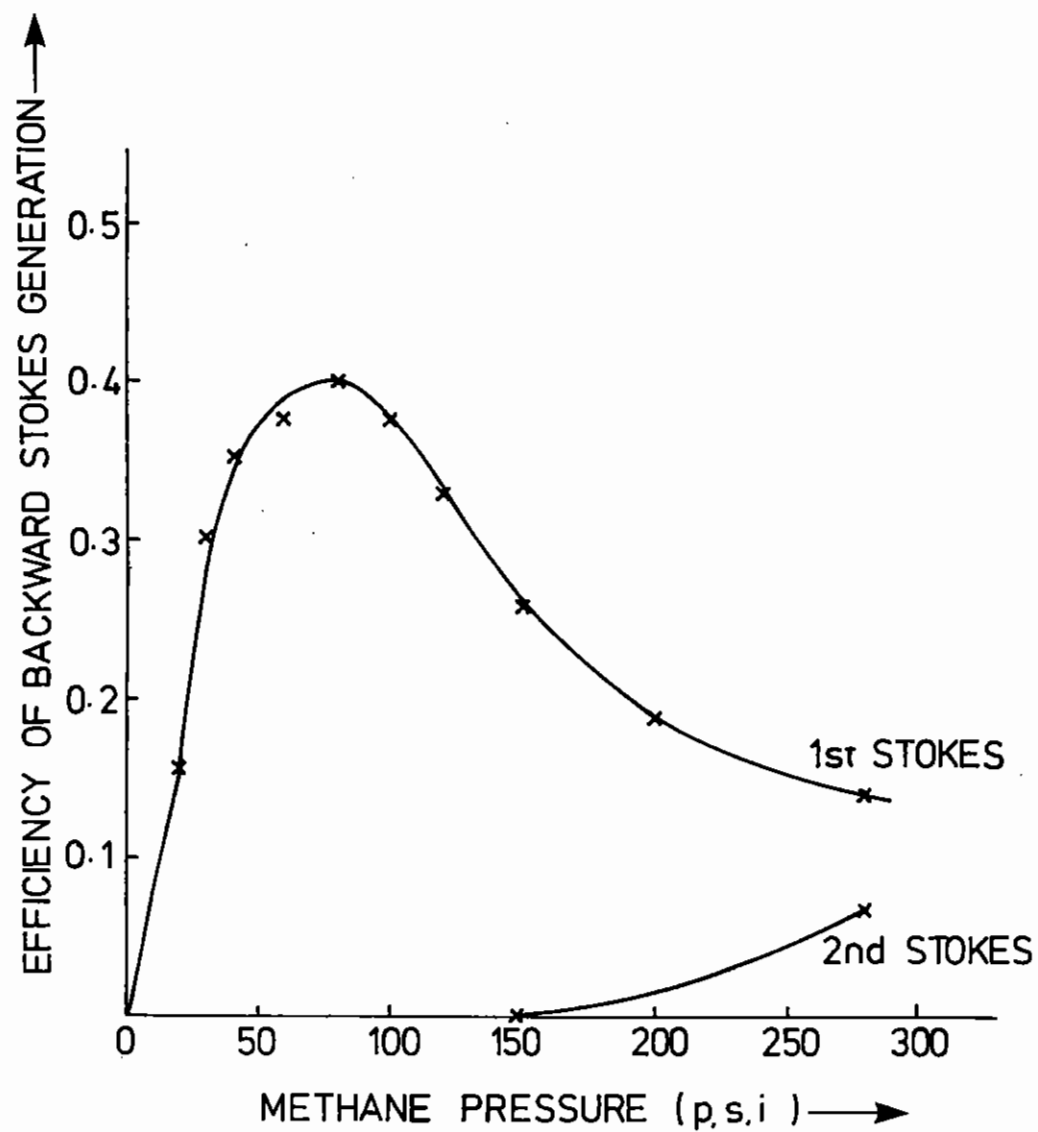


Fig. C2.9 Efficiency versus methane pressure curves for backward generated 1st and 2nd Stokes.

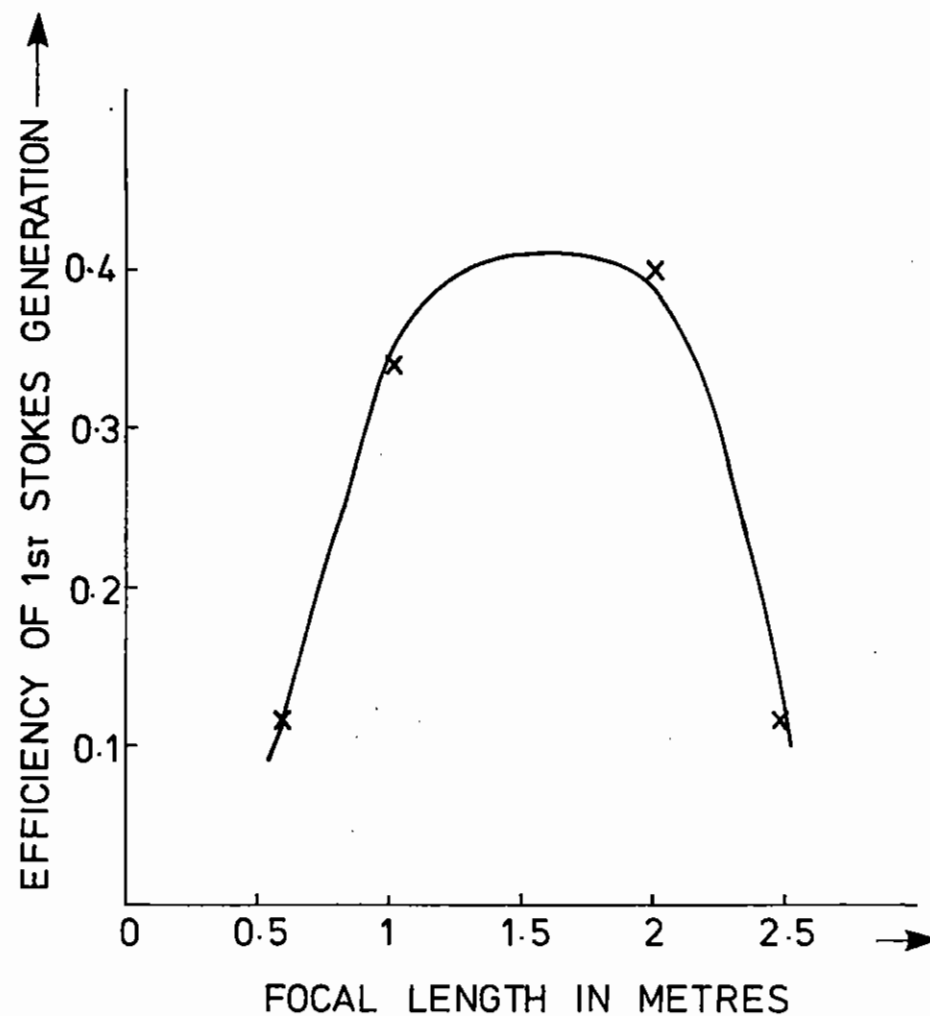
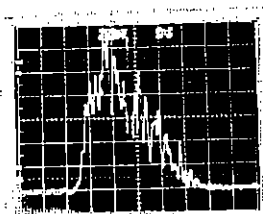


Fig. C2.10 Efficiency as a function of the focal length of the lens used to focus the KrF laser into the methane cell.

(A)

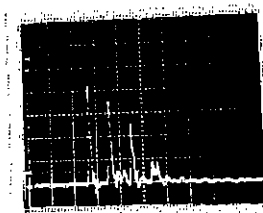
PUMP



5ns/div

(B)

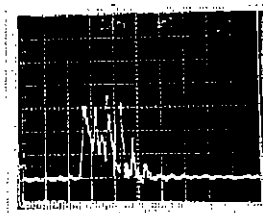
1st STOKES
WITH LONG
FOCUS LENS
($f = 2.5\text{m}$)



5ns/div

(C)

1st STOKES
WITH SHORT
FOCUS LENS
($f = 1\text{m}$)



5ns/div

(D)

2nd STOKES
AT HIGH PRESSURE
($f = 2\text{m}$)



0.5ns/div

bandwidth of the Raman process. At the first Stokes the time evolution is less spiky particularly for a shorter focal length lens for which the Raman gain length is less. However peak powers for this component are still comparable to that of the pump. The presence of heavy modulation on the pump pulse on a nanosecond timescale will strongly encourage short pulse generation in the backward Raman components. A more temporally uniform Stokes beam may be achieved when the preamplifier is far more heavily saturated than in this experiment (which uses it only in single pass). The temporal modulation on the pump beam will be suppressed by this nonlinear saturation process.

Fig. C2.11

Oscilloscope of:

- (A) KrF pump pulse into generator
- (B) Backward 1st Stokes. $f = 2.5\text{m}$ lens.
- (C) Backward 1st Stokes. $f = 1.0\text{m}$ lens.
- (D) Backward 2nd Stokes. $f = 2.0\text{m}$ lens.

References

- C2.1 I V Tomov, R Fedosejevs, D C D McKen, C Dornier and A A Offenberger, Opt Letts, 9, 8 (1983)

C3 UV LASER DAMAGE STUDIES

D C Emmony, I Laidler, L A Mahmood and S S Wiseall (Loughborough)

C3.1 Introduction

Laser damage studies in the past year have been concerned with improving diagnostics, laser damage in aluminium mirrors (C3.1) and an investigation of laser damage threshold definitions. An absorbing impurity model has been suggested which is able to predict damage thresholds and damage probabilities (C3.2).

Exploratory work has begun on an acoustic probe technique to detect the early stages of the onset of damage.

A CASE project in collaboration with BDH Chemicals Limited, has recently begun. This project involves the study of optically induced primary defects in single crystal CaF_2 and other UV materials.

A variable length liquid attenuator has been developed which allows uniform attenuation of the excimer laser at 249 nm. Cobalt and nickel sulphate or Rhodamine 6G laser dye have been found to be suitable absorbers. The use of an attenuator has two benefits. Laser damage measurements can be carried out in the same focal region and focused excimer beam profile measurements are possible using a pinhole and small area pyro-electric detector.

C3.2 UV Laser Damage in Aluminium Mirrors

An investigation of KrF laser damage processes in aluminium mirrors has been conducted (C3.1). Optimum coating deposition conditions can produce mirrors with a reflectivity greater than 90% at 249 nm. The aluminium oxide layer which grows naturally (~ 4 nm) does not affect the reflectivity of the mirror at 249 nm. Laser induced heating of the aluminium coating appears to be the damage limiting process. Damage tests were also performed on silicon oxide, ThF_4 and MgF_2 overcoated aluminium mirrors. There was no large variation in thresholds between coatings, being typically in the range 0.2 - 0.4 Joules cm^{-2} .

SILICON OXIDE OVERCOATED ALUMINIUM MIRROR

The damage limiting process for MgF_2 and ThF_4 overcoated mirrors was the adherence of the coating to the aluminium. Some interesting observations were made on silicon-oxide overcoated mirrors. The standard silicon oxide overcoating on commercial mirrors is the non-stoichiometric oxide, SiO_x , where x lies between one and two. UV irradiation with a Mercury discharge was found to dramatically increase the UV reflectivity, Figure C3.1. The improvement increases with irradiation time and is larger for shorter wavelengths. Ellipsometry measurements at 633 nm suggested that the film thickness and refractive index at this wavelength were not influenced by the irradiation. ESCA was employed to measure the stoichiometry of the SiO_x film as a function of irradiation time. This analysis showed the surface layer to be fully oxidised, with a stoichiometry index of 2. For UV irradiation times of less than 90 minutes there is a gradual increase in the stoichiometry value with sputter depth and irradiation time from about 1.5 to 1.75. Irradiation times of 150 minutes produce a fully oxidised coating, SiO_2 . The largest improvement in reflectivity occurs for irradiation times of less than 40 minutes, in this case the stoichiometry value changes very little. Similar observations have been made by Hass et al (C3.3) who attributed part of the increase in reflectivity to an improved structural ordering of the SiO_x overcoat layer.

C3.3 Probability Model

A model has been examined which considers the most damage sensitive impurities to be distributed with a Poisson distribution at thin film interfaces. A theoretical model considering pulsed laser induced heating of a lightly absorbing spherical particle in an infinite dielectric medium has also been considered. The energy density required to heat the centre of the absorbing particle to its melting point is shown in Figure C3.2. The absorbing particle is assumed to be silicon oxide (SiO_x) in a non-absorbing silicon dioxide host medium. The model assumes uniform heating throughout the particle. The imaginary part of the complex refractive index of the absorbing particle limits the maximum particle radius for which the assumption of uniform heating is valid. In Figure C3.2 for $n_i = 0.001$ the maximum radius is about $0.9 \mu\text{m}$. The predicted minimum energy density of 8.3 J cm^{-2} for particle radii $\sim 0.5 - 1.0 \mu\text{m}$ is 2 - 4 times higher than that

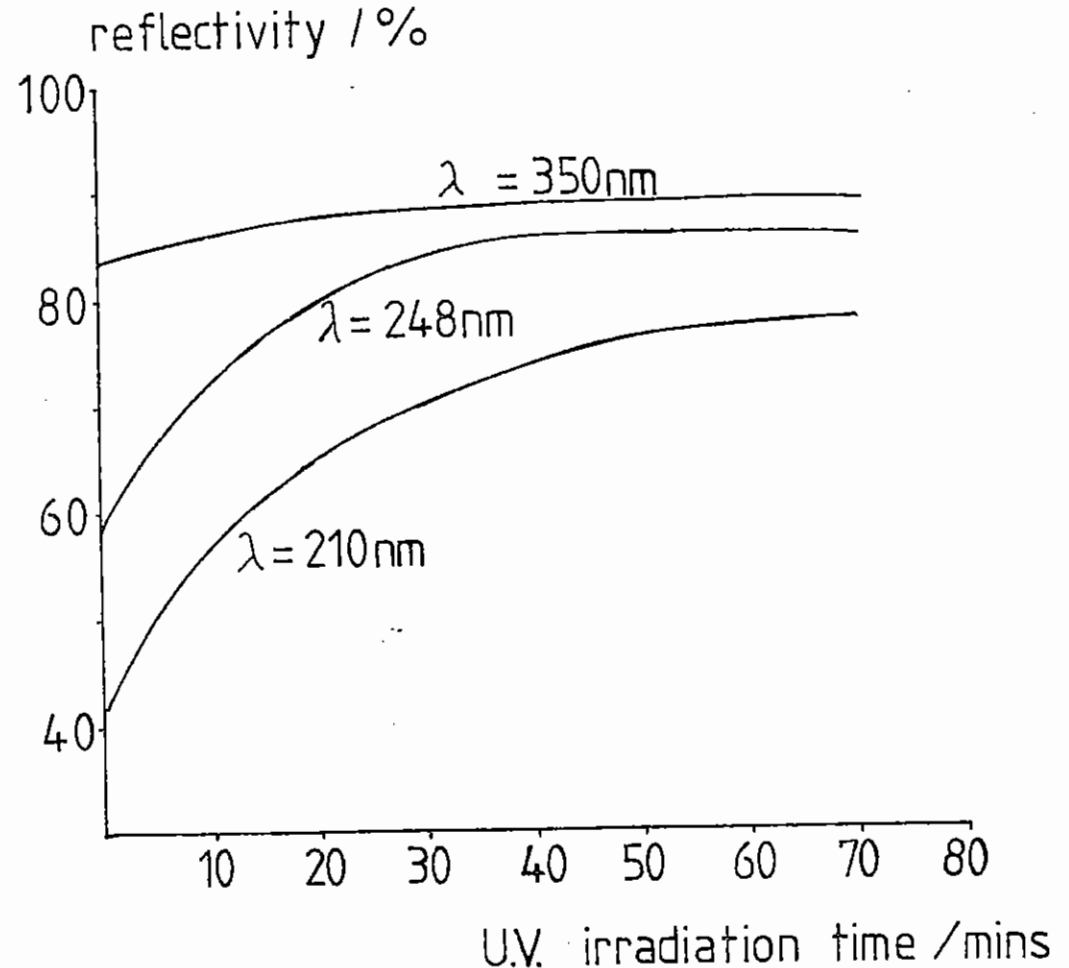


Fig. C3.1 Effect of UV irradiation on mirror reflectivity.

energy density
/ Jcm⁻²

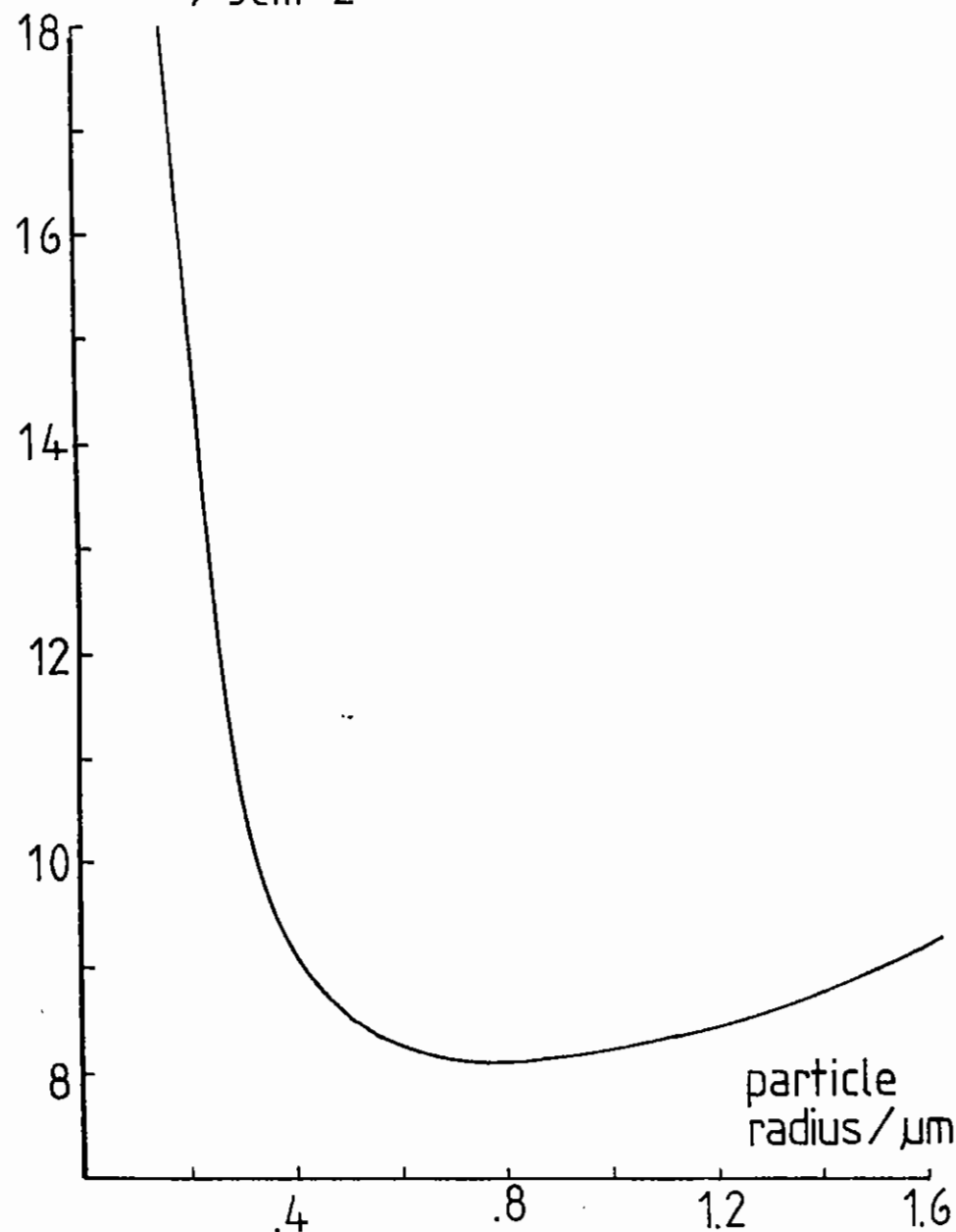


Fig. C3.2 Theoretical energy density required to raise the centre of an absorbing impurity to its melting point.

observed experimentally. For such small particles the absorption cross section is given by Mie theory and is $\ll \pi a^2$, its geometrical value. Guenther (C3.4) has found that experimental data scales with the geometrical cross section even for $\lambda/2$ films for which the impurity size would be expected to be $< \lambda/2$. This would be expected to occur for moderately absorbing particles with $n_i > 0.01$. The requirement of uniform heating throughout the particle, limits the particle size to $< 0.1 \mu\text{m}$. The analysis for larger particles requires a spatially dependent source term in the heat conduction equation.

The results of the damage probability model are illustrated in Figures C3.3. They are expressed as a damage probability versus peak test fluence. The test specimen was a three layer non-quarter wave antireflection coating designed for use at 249 nm. One half of the sample had been irradiated for 30 hours with a Mercury discharge lamp. The experimental points were determined by damage testing the sample at 10 to 20 sites at each fluence and noting the fraction that damaged. The theoretical lines were fitted using the defect density as a fitting parameter. The results suggest that pre-irradiation with a UV discharge lamp source does not alter the damage threshold of the most damage sensitive defects but decreases the defect density from 1.7 mm^{-2} to 1.0 mm^{-2} .

The gradient of the theoretical line is controlled by the defect density and laser spot size (C3.5). Laser damage thresholds are defined in several different ways but experimentally they converge for large beam area tests.

Some interesting accumulative processes have also recently been observed. Figures C3.4 shows the shot number required to produce ejection of material from the test specimen plotted against peak fluence. The results may be explained by an accumulative increase in the energy absorption coefficient. The increment in absorption coefficient per shot is less for the irradiated section. More shots are therefore required to produce damage. If the increment in the absorption coefficient/shot was constant the product of shot number and fluence in Figure C3.4 would also be constant. This is not the case and a non-linear increase in the absorption coefficient with irradiation level occurs. The accumulative process appears only to occur

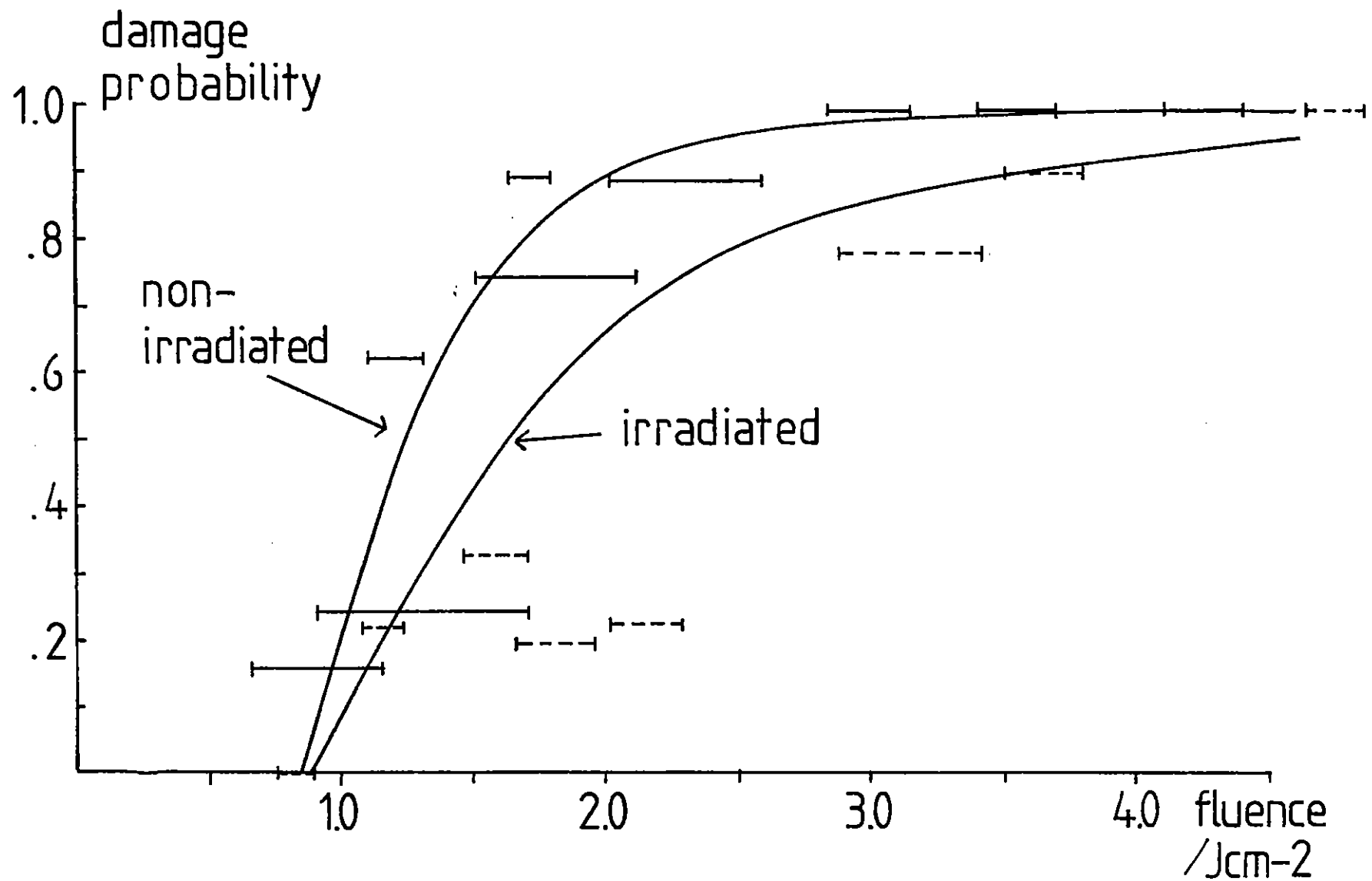


Fig. C3.3 Damage probability as a function of laser fluence.

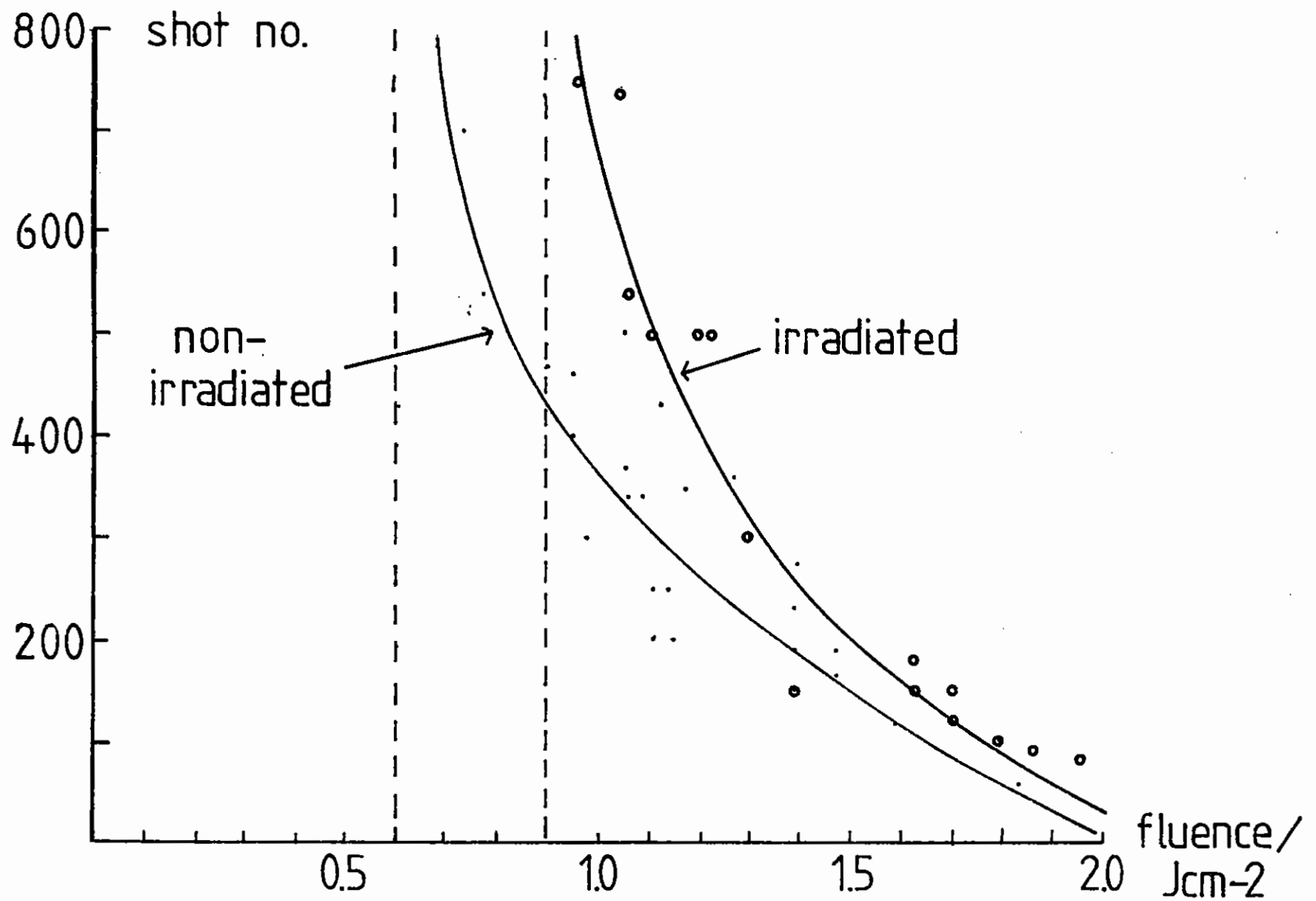


Fig. C3.4 The number of laser shots required to produce visible damage as a function of laser fluence.

above certain threshold energy densities.

C3.4 Damage in Synthetic Crystals

The CASE project in collaboration with BDH Chemicals Limited, is in its earliest stages. Previous work has shown that CaF_2 and other UV transmitting materials fluoresce and develop colour centres in excimer radiation. The project seeks to understand the mechanisms and the role of laser induced defects on the breakdown threshold of CaF_2 in the first instance. Absorption and fluorescence spectroscopy are being used to study the interaction process. A variable delay pulse derived from a simultaneously pumped tunable dye laser is to be used to probe the damage threshold. It is hoped to isolate the competing kinetic processes involved in optical breakdown while at the same time monitoring impurity species in the material.

C3.5 Photoacoustic Detection

The onset of damage and definition of damage thresholds has received much attention. A simple acoustic technique has been developed which may overcome some of the difficulties and which lends itself to spot size dependence studies. A piezoelectric transducer is mounted directly on the component to be damage tested and the amplitude of the acoustic transient recorded as a function of energy density. Figure C3.5 shows the acoustic amplitude as a function of CO_2 laser pulse energy. At low energy densities the acoustic signal is due to a thermoelastic mechanism and is linearly dependent upon the pulse energy. This dependence changes discontinuously at the onset of damage. The technique shows excellent agreement with damage thresholds measured in other ways on a wide range of substrates and mirrors including Al, stainless steel, Cu and Mo, as well as a number of high reflection and anti-reflection thin film systems. So far the technique has only been used as $10.6 \mu\text{m}$ but it is planned to extend this work to the detection of the onset of damage in the ultra violet.

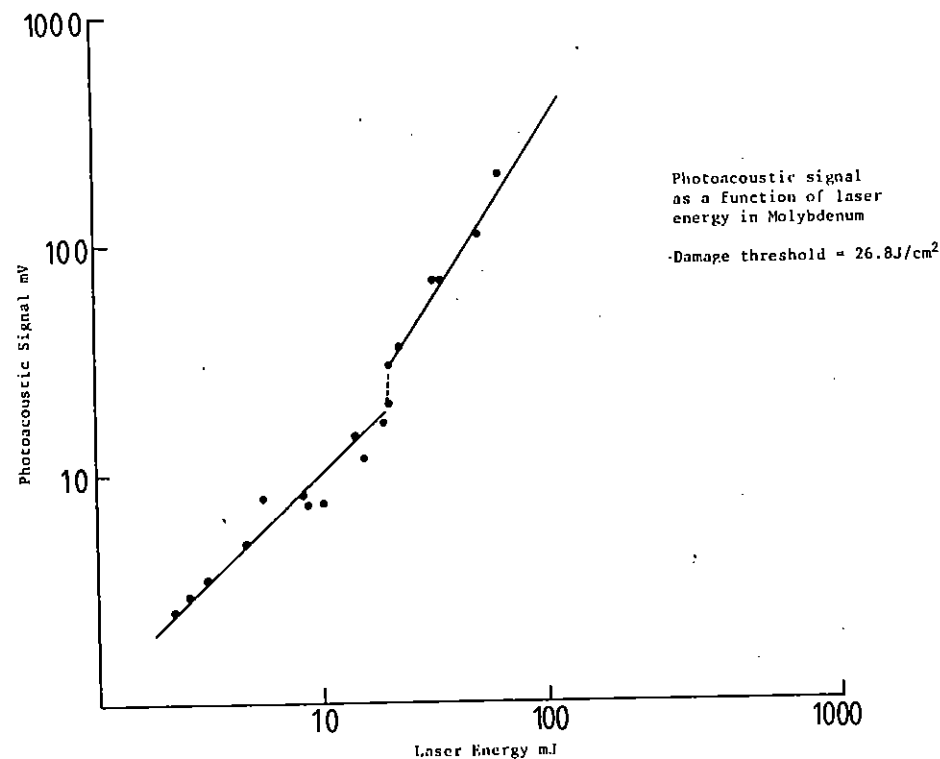


Fig. C3.5 Acoustic signal in Molybdenum with CO_2 laser energy showing the onset of damage.

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C4

E-BEAM PUMPED KrF LASER SYSTEMS

C4.1

Sprite Laser Studies

C B Edwards, F O'Neill, M J Shaw (RAL), Z Osborne and C Phipps (LANL)

The Sprite laser system is now fully operational and during the reporting year approximately 600 laser shots were fired mainly in connection with the Raman amplifier experiments described in section C5 below. The laser system was shutdown late in 1983 for a period of four months for a fairly extensive re-build of sections of the pulsed power equipment and for the installation of a new laser cell, the old one having developed a large gas leak. Our experience with glued pressure foils in the laser cells has been good with the glue joints proving to be extremely robust in operation and well capable of surviving the large (100-200 torr) pressure pulse generated when the e-beams are dumped into the laser gas. The old laser cell had been in Sprite for about one year and had deteriorated due to F_2 attack of the epoxy glue. In the new cell just installed only a very thin layer of glue is exposed to the F_2 and degradation is not expected to be a problem.

One of the main problems encountered in operating Sprite has been with F_2 attack of the laser optical coatings which are used internally in the cell. The usual laser cavity consists of an unstable resonator with a 260 mm aperture, 4 m radius of curvature rear 100% reflector in contact with laser gas and the output end of the cell closed off with a 260 mm aperture fused silica AR/AR coated window. An externally mounted convex mirror, 65 mm diameter, 1 m radius of curvature completes the resonator. The large rear mirror is an all-fluoride multilayer dielectric coating on a pyrex substrate and we find that this withstands F_2 attack over long periods provided exposure of the coating to the laboratory atmosphere is minimised. These coatings have been measured by Loughborough University to have a laser damage threshold $\sim 2 \text{ J/cm}^2$ when new which drops to $< 1 \text{ J/cm}^2$ after F_2 exposure. In any case the damage threshold is adequately high for use at typical KrF laser intracavity fluences of $< 500 \text{ mJ/cm}^2$.

Experience with AR coatings in contact with F_2 has not been as good. The

C.16

C.15

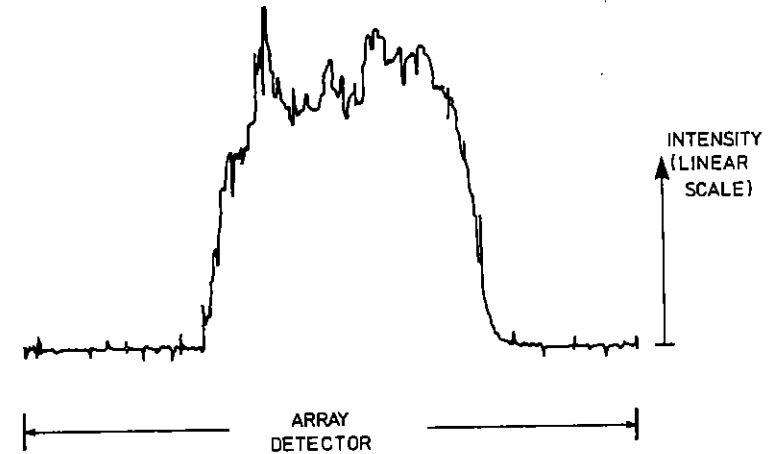
SPRITE LASER BEAM PROFILE

coatings are typically of two-layer design, one layer being a fluoride and the other an oxide. It seems that F_2 (or HF) rapidly attacks the oxide layer inducing a cloudy appearance and making the layer easily damaged by the laser. More work will have to be done with AR coatings of different designs to improve their F_2 compatability.

The Sprite laser has been mainly operated with an unstable resonator cavity the aim being to generate low divergence beams of good spatial quality. The laser beam uniformity has been measured using a 128 element pyroelectric detector array with 1 mm centre-to-centre spacing. Measurements have been taken with the array mounted in a vertical plane and in a plane at 45° to the vertical to see if the four sided Sprite e-beam geometry is in anyway imprinted on the laser beam. Figures C4.1(a) and (b) are two such traces of the laser beam intensity profile showing a peak to valley uniformity of $\pm 30\%$ in both planes. For these measurements $\sim 8\%$ of the Sprite beam was picked off by a beam-splitter and compressed onto the array using a 1.2 m focal length lens. This beam uniformity is sufficiently good for pumping Raman amplifiers and is adequate for most target experiments. The array trace (a) corresponds to a plane which passes between e-beam diodes and the low laser intensity at the edges is as a result of reduced e-beam pumping in these regions. The 45° trace (b) corresponds to a plane passing through two e-beam diodes and have the enhanced pumping close to the e-beam foils results in the laser beam profile peaking at the edges. Nevertheless the beam profile in both planes is quite uniform showing that the e-beam scattering in the laser gases effectively "washes-out" the four-sided pumping geometry.

An important requirement of Sprite for pumping Raman lasers was that it should operate with low beam divergence and narrow line-width ($< 0.3 \text{ cm}^{-1}$). The method chosen to achieve this was to operate with an $M = 4$ unstable resonator cavity and to injection lock it with a pulse from a low divergence, line narrowed Lambda Physik EMG150 laser. The optical set-up is shown in Figures C4.2. The injection pulse was carefully timed to enter the Sprite cavity just as the gain in the large laser was starting to rise as a result of the e-beam excitation. The effect of Sprite injection was very dramatic resulting in the following:

(a)



(b)

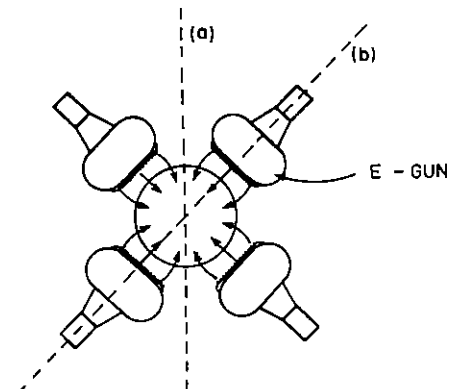
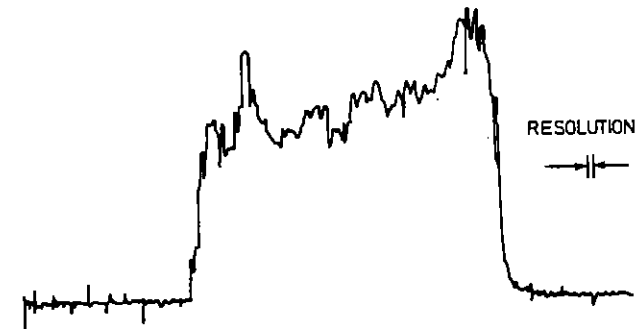


Fig. C4.1

Sprite laser beam profiles taken by compressing a fraction of the output onto a pyroelectric array using a long focus lens ($f \sim 1.0\text{m}$). Profiles (a) and (b) are taken along two different beam diameters.

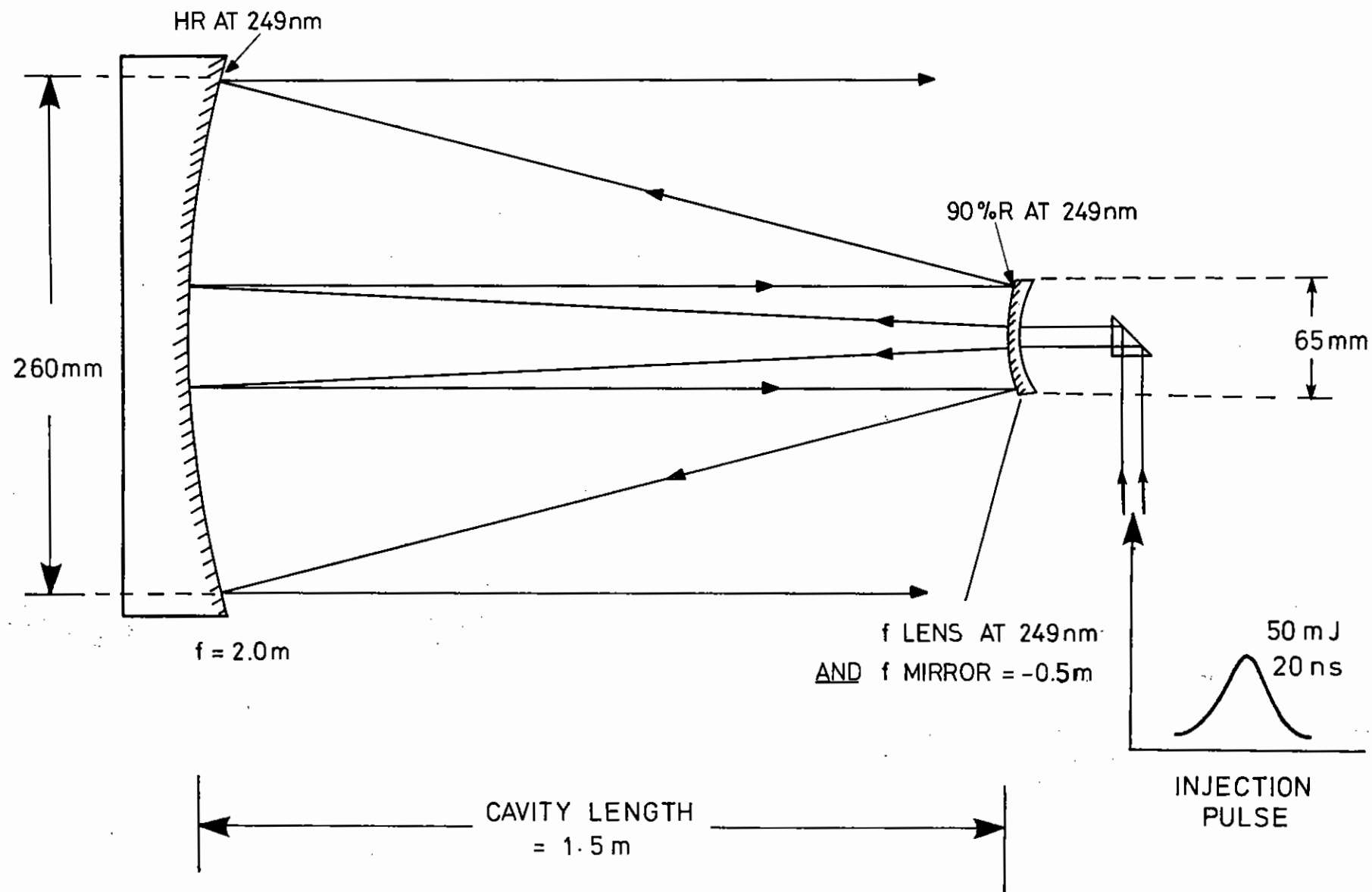


Fig. C4.2 Sprite injection locked unstable resonator cavity.

SPRITE INJECTION LOCKED PERFORMANCE

- (i) Appearance of high contrast diffraction fringes on the far field beam pattern due to beam clipping by hard edges in the cavity. This is evidence of the low beam divergence - in the non-injected beam these diffraction fringes are washed out in the far field because of the higher beam divergence.
- (ii) Increased laser pulse-width and output energy. Figure C4.3 shows the laser pulse width, with and without injection where the corresponding laser energies were 142 J and 120 J.
- (iii) Line-narrowed output with a line-width of $< 0.3 \text{ cm}^{-1}$ (see Figure C4.3). As yet we have not measured the percentage of the laser energy in the narrow line but from interferometer measurements it appears to be very high ($> 80\%$).

In general therefore our operating experience with Sprite has been good. Our aim has not been to push the machine to its limits and we have typically operated with laser energies in the range 100 - 150 J which are easily achieved at modest running voltages. These laser pulse energies are obtained using unstable resonator cavities and for a deposited e-beam energies $\sim 2500 \text{ J}$. The maximum pump energy that can easily be achieved is $\sim 3000 \text{ J}$ and by using plane/plane resonator optics this should give an energy output $\sim 300 \text{ J}$ assuming 10% laser efficiency. This 300 J is also the level of energy that should be available when Sprite is used as a saturated amplifier in the KrF multiplexer arrangement.

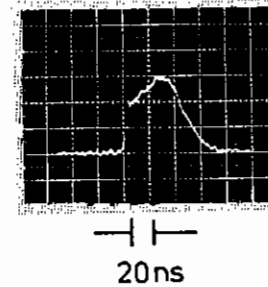
C4.2 Sprite Pulsed Power Performance

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

The major change in the Sprite pulsed power system in the past year has been the introduction of a new triggering system which includes laser triggering of the main water capacitor SF_6 switch for the first time. This modification was necessitated by the introduction of the Lambda Physik oscillator to the Sprite system. This laser's minimum trigger delay of

INJECTION
ON
(142 J)

(a)



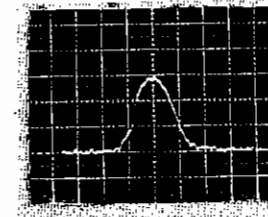
(b)



$F \text{ SR} = 1.5 \text{ cm}^{-1}$

INJECTION
OFF
(120 J)

(c)



(d)



Fig. C4.3 Sprite injection locked performance. Laser pulse shapes and Fabry Perot interferograms with and without injection.

around 600ns meant that a trigger pulse had to be provided before the water capacitor switch fired - the previous timing point of the Sprite system. The solution adopted is shown in Figure C4.4 which shows the layout of trigger pulses to the various parts of the system. The main trigger pulse is derived from the water capacitor voltage monitor rather than the Marx trigger generator thus avoiding the jitter inherent in the Marx erection. A level detector provides a pulse close to the beginning of the water capacitor charge cycle. The pulse is used to fire a CX1622 Hydrogen thyratron which provides a high level (30 kV) pulse via a delay line to the first trigger laser and a low voltage (15 V) pulse directly to the Lambda Physik oscillator laser. This is thus triggered before the first trigger laser, the output of which is used to trigger the water capacitor spark gap. A second delay line takes the high voltage pulse to the second trigger laser which is used to fire the four pulse forming line gaps.

Since the jitter in the laser triggered gaps can be very small ($\sim 1 - 3$ ns), the overall system can be made to perform under accurately repeatable conditions. The delay lines can be adjusted to give optimal triggering of the water capacitor and diodes. When so adjusted, over 90% of the energy stored on the water capacitor is delivered to the diodes within the useful pulse duration. The main inefficiency in the system is now the diode/gas interface where typically only 40-50% of the diode energy gets into the gas. This is to be compared with 50-60% measured previously with a test diode under similar conditions. The difference is most likely due to losses on the curved main foil and at the curved laser chamber wall which gives a larger surface loss than on the single flat face of the test diode (see 1982 Annual Report).

Operating experience over the past year has provided useful data on the system. For example, laser triggering has proven to be highly reliable the only problem being the etching of the fused silica focusing lens by (presumably HF impurities) formed in the SF_6 gas after breakdown. This problem has been minimised by mounting the lens in a separate chamber fed with fresh SF_6 and communicating to the gap only via a small hole in one electrode. The four PFL gaps were dismantled after approximately 500 shots for cleaning and inspection. They were found to be in excellent condition

with no evidence of surface tracking or excessive sputtering.

On the other hand serious problems have been found in the 4" diameter copper feed lines to the diodes. Firstly, after a period of about a year it was found that the inside of the copper lines were becoming plated with a substantial layer of black CuO . Since this is an insulator the effect is to reduce the breakdown threshold in this part of the system. It is not entirely clear what the mechanism producing the oxide layer but it is probably caused by H_2O hydrolysis products, produced in streamers, reacting with the copper. Copper should certainly be avoided in future water filled pulsed power systems. The bronze parts of the lines (bends and tees) seem unaffected by this problem. Secondly, breakdown in this part of the system has been occurring far more often than originally envisaged. This is almost certainly due to the small residual voltage which remains on the lines after firing. Since there is very little to dissipate this charge after the main pulse is over it finds the weakest point which is in the small diameter output lines. These breakdowns give rise to severe pressure pulses which have caused bulging of the lines and failure of soft soldered joints. The elimination of soft soldered joints and the provision of extra bellows sections have been tried to reduce the effects of the breakdowns but the root cause remains. Tin plating is being tried in a section of the system to hopefully overcome the CuO problem, but some redesign of this section of the system to (a) reduce the fields and (b) provide some form of energy dissipation will be necessary before the full potential of the Sprite pulsed power system can be realised on a routine basis.

C4.3 Goblin Pre-Amplifier Progress

D Baker, D Craddock, E Madraszek, H T Medhurst and M J Shaw (RAL)

A major new component of the multiplexed Sprite system is the e-beam pumped KrF laser pre-amplifier, the function of which is to boost the output of the oscillator to a level sufficiently high that the main Sprite amplifier is efficiently extracted. This system has acquired the code name Goblin and uses components in common with both ELF and Sprite. The nominal output of Goblin is 20 J in a 100 ns pulse from a 75 mm diameter aperture. Initially

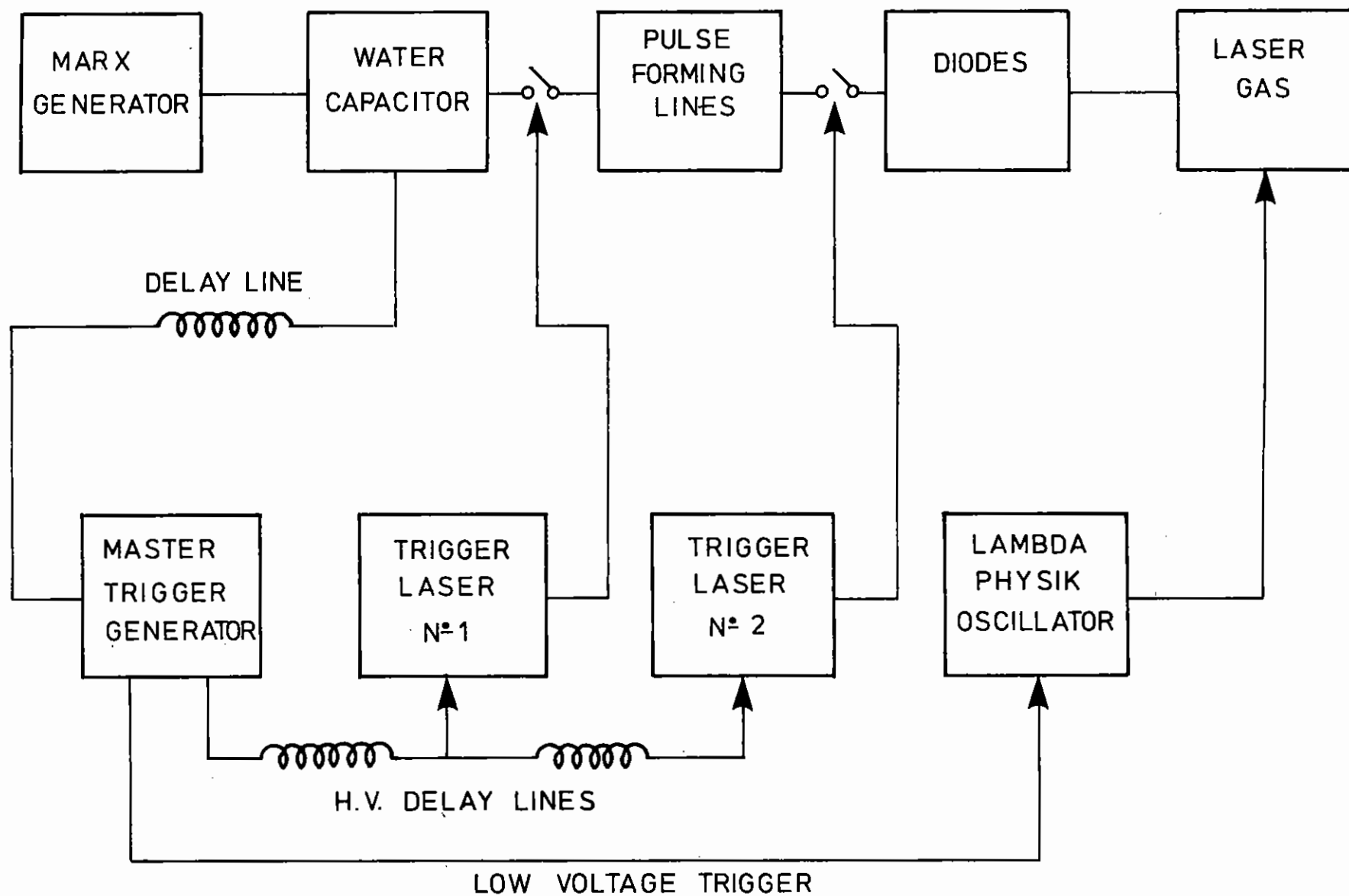


Fig. C4.4 Block diagram of Sprite pulsed power and triggering system.

it is proposed to use the old ELF diode chamber and laser cell. This is fed by a Sprite type pulse forming line with a laser triggered output switch. In this case the PFL is 2 m long giving a pulse duration of 120 ns. The line impedance is 5 Ω and the 12 nF line stores 3.8 kJ at its maximum charge voltage of 800 kV.

An entirely new Marx generator has been designed to power Goblin. It comprises 10 stages of 0.22 μ F (Maxwell S series) capacitors and stores a maximum of 6.2 kJ at 75 kV charge voltage. The voltage gain from marx to line is calculated to be 1.1 and the energy transfer efficiency 70%. The charge time of the pulse forming line will be approximately 400 ns. The most important parts of the Marx generator are the switches and it was decided to use the same switching arrangement as in ELF where midplane triggered ball gaps are arranged in a vertical column in a common tubular pressure vessel. Figure C4.5 shows the Goblin Marx switch column during assembly. The experience with the ELF switch column has been very good - several thousand shots having been obtained without significant degradation in performance.

C4.3.1 Marx Generator Triggering

One new requirement for the Goblin Marx generator (and also as a retrofit for the ELF Marx) is that of low jitter triggering. This has previously not been a problem since everything could be slaved to the water capacitor charge cycle which occurs after the Marx is triggered. For the Sprite multiplexer however it will be required to trigger both the ELF and Goblin Marx generators in correct time sequence with a total jitter of 25 ns or so.

A triggering system requiring minimal attention and having very high reliability with low prefire rate and capable of low jitter operation from a low voltage timing pulse was required. These requirements have led to a fairly novel design employing a hydrogen thyatron and a magnetic switch. The basic circuit is shown in Figure C4.6. Energy is stored on C_1 and switched by the tetrode thyatron V_1 . V_1 and C_1 sit outside the Marx tank and are connected to C_2 , L_2 and T_1 under oil in the Marx tank via a cable

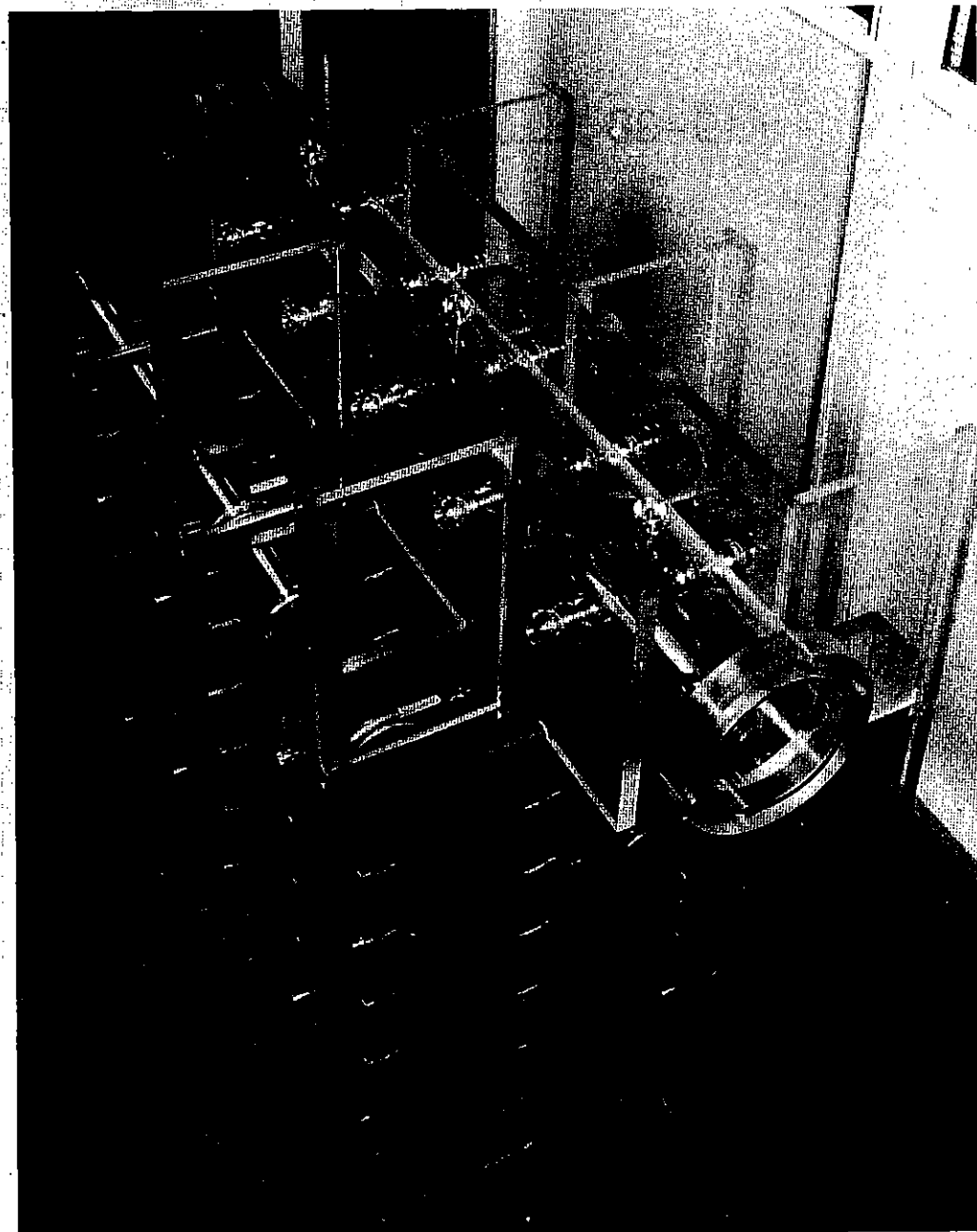


Fig. C4.5 Goblin Marx column during assembly.

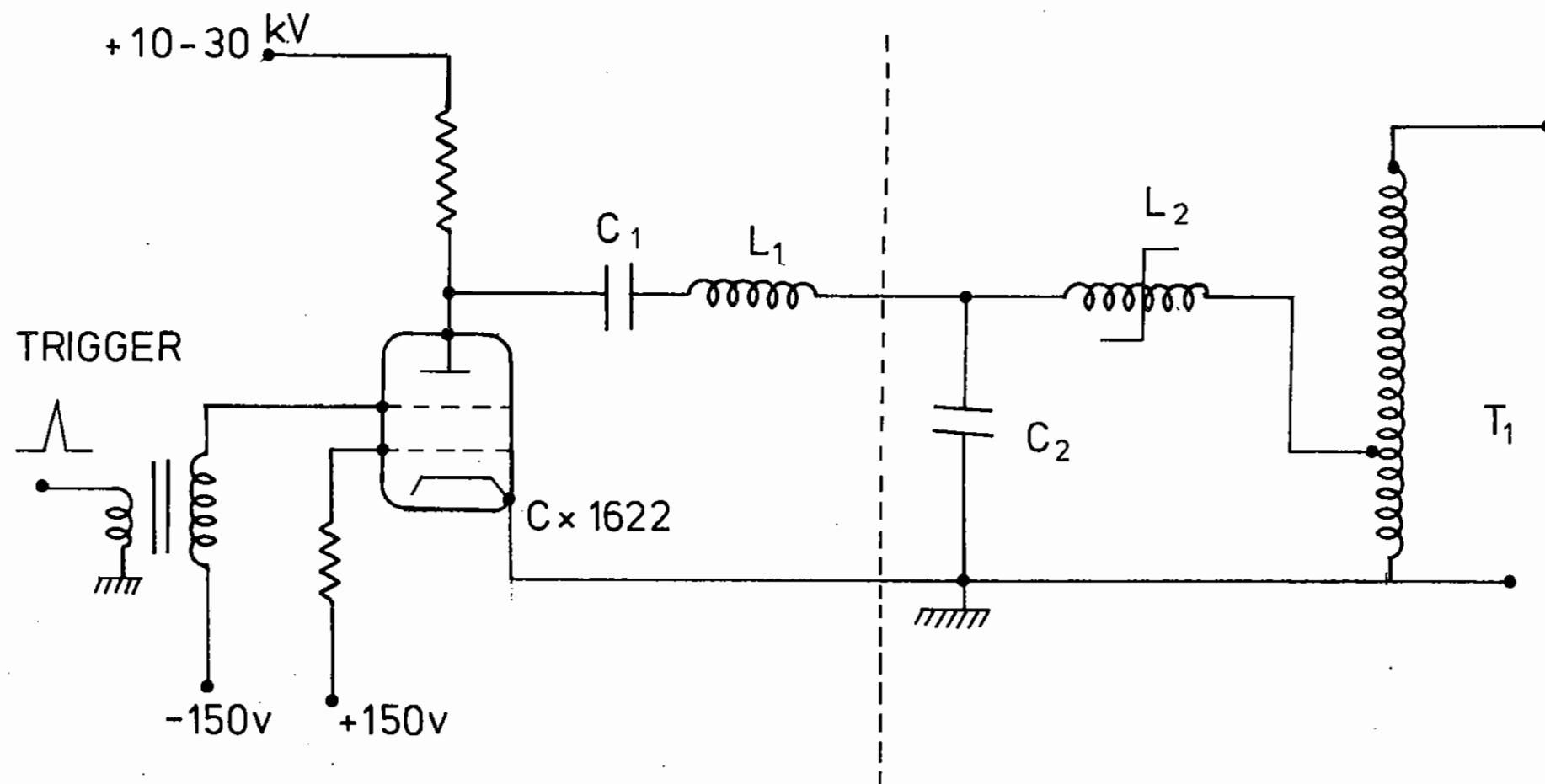


Fig. C4.6 Marx trigger generator utilising a magnetic switch.

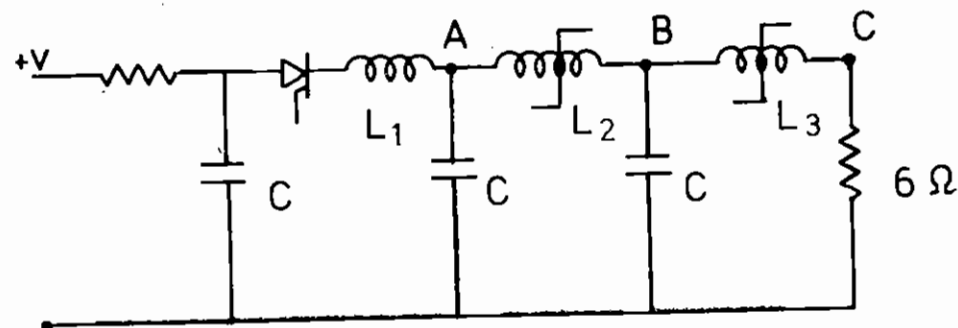
which contributes most of the inductance L_1 . This inductance is chosen such that the peak current and peak di/dt are well within the thyatron ratings. L_2 is a saturable inductor so chosen that L_2 (unsat) is $\gg L_1$ and L_2 (sat) is $\ll L_1$. Under these conditions C_2 charges until

$$\int v(C_2)dt = N A \Delta B$$

where $N A \Delta B$ are the number of turns, cross sectional area and saturation flux of L_2 respectively. At this precise point L_2 switches very rapidly from its unsaturated to its saturated state allowing current to flow to the load. The current risetime is determined solely by the saturated inductance of L_2 and can be many times (typically 5 or so) shorter than the risetime at the thyatron. T_1 is an air-cored high voltage pulse transformer the output of which is connected directly to the Marx trigger electrodes. Typical output voltage required for low jitter triggering is ~ 300 kV with a risetime of < 100 ns. The advantage of this scheme is that the thyatron is operated under relatively benign conditions so that long life and high reliability may be expected. Also the components in the Marx oil tank are purely passive in nature and, providing electric fields in T_1 are considerably less than the breakdown field, should be maintenance free. In order to test out the idea a small scale version using an SCR switch and ferrite cores was set up as shown in Figure C4.7. This two stage system clearly shows the decrease in the risetime and hence current gain per stage. Ferrite is not particularly useful for high power applications since as has been shown by Birx (C4.1) the energy handled per stage is given by:

$$E = \frac{4V \Delta B^2}{G^2 \pi^2 \mu_s}$$

where V is the core volume, ΔB the saturation flux in Tesla, G the stage gain and μ_s the saturated permeability of the core (generally close to μ_0). For ferrite ΔB is rather small (typically 0.3 T) giving rise to specific energies of about 2 J/l for a gain of 5. On the other hand, a new magnetic material, Metglas, has high permeability and high frequency response as well as a much higher saturation flux (~ 1.5 T) which gives specific energies ~ 50 J/l. A switch using Metglas is at present under construction for use in this trigger generator application.



$C = 1.0 \mu F$
 $L_1 = 35 \mu H$ air cored
 $L_2 = 2.1 \mu H$ (unsat)
 $L_3 = 120 \mu H$ (unsat)

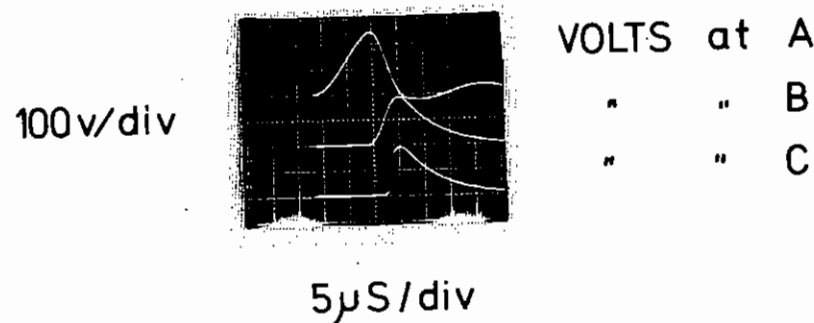


Fig. C4.7 Magnetic switch test circuit.

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C5 LIGHT GUIDED FORWARD RAMAN AMPLIFIER EXPERIMENT

C B Edwards, F O'Neill, M J Shaw, G Levy and Y Matumoto (RAL)

C5.1 Experimental Results to Date

The basic Raman amplifier experiment was described in detail in last year's annual report. The main details of the experiment are shown schematically in Figure C5.1 and consists of a 1.8 m long square section light guide constructed of float glass plates giving a $2.5 \times 2.5 \text{ cm}^2$ aperture. The guide is pumped by 56 plane beams which are split from the Sprite output using a multi-element segmented mirror. The beams are plane so as not to bring the beams to a focus in the Raman cell which could result in gas breakdown and spontaneous Raman generation. The large number of beams were dictated by the need to achieve a sufficiently high intensity in the guide using the 60 ns Sprite pump pulse.

Alignment of the segmented mirror (after trying a number of poorer methods) was best achieved by using the oscillator laser which is used to inject Sprite and which, after four passes in the Sprite cavity is expanded to full beam size. A system of beam stops is used which illuminates each element of the segmented mirror in turn. Each element is then adjusted so that its beam is centred on the light guide entrance. The accuracy of alignment using this technique may be judged from Figure C5.2 which shows a burn mark on a silicon wafer placed at the light guide entrance. Each mirror segment was designed to produce a $23 \times 23 \text{ mm}^2$ spot to allow for $\pm 1 \text{ mm}$ misalignment at the entrance of the $25 \times 25 \text{ mm}^2$ guide. The burn mark in Figure C5.2 which comes from the overlay of all 56 beams is in fact exactly $23 \times 23 \text{ mm}^2$ indicating that alignment is well within tolerance.

The energy available at the guide entrance was somewhat less than anticipated. The maximum being about 40 J in a 50 ns pulse. The resulting average intensity in the guide was thus 130 MW cm^{-2} . The guide transmission was measured using joulemeters at the entrance and exit apertures and found to be 70% - somewhat less than the calculated 85%. The main experiment consisted of measuring the energy of the Stokes beam at the output of the

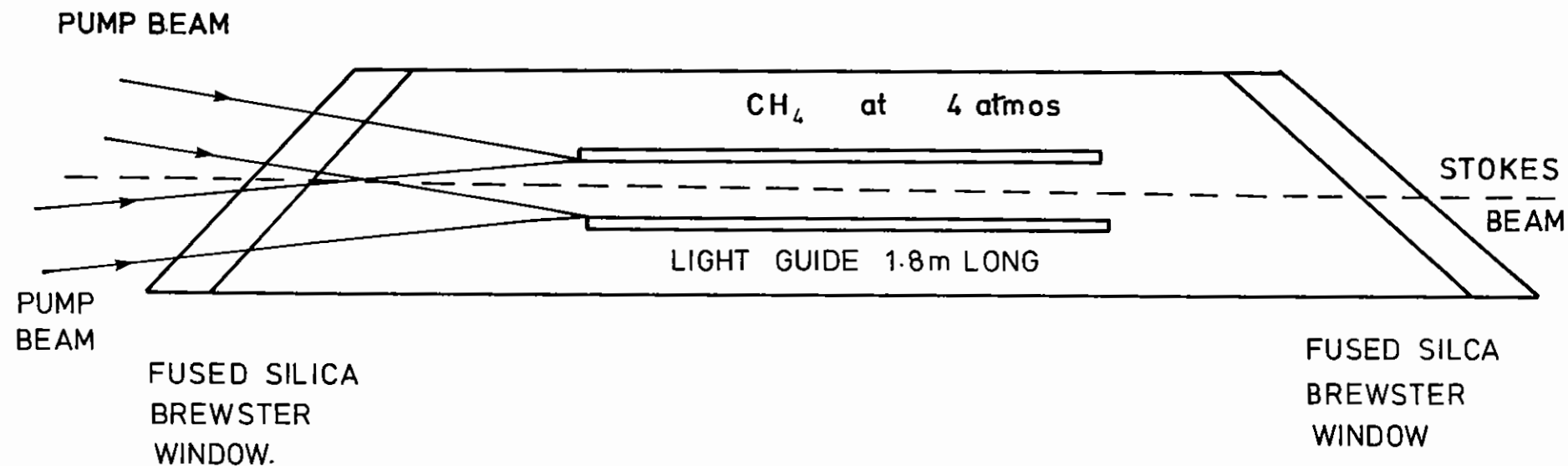


Fig. C5.1 Schematic diagram of the light guided Raman amplifier.

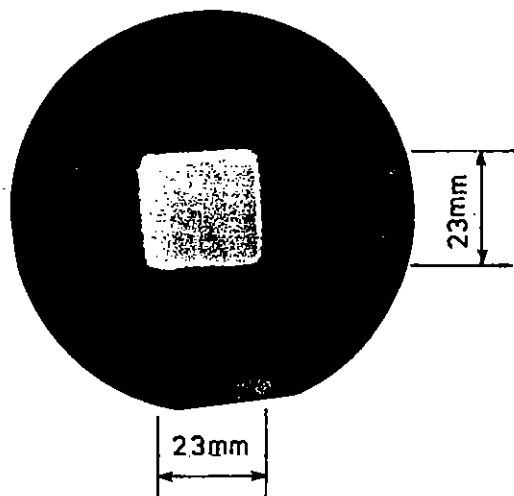


Fig. C5.2 Laser burn pattern at entrance to light guide with all 56 beams from the segmented mirror overlaid.

guide with and without pumping. The fact that the pump beams did not co-propagate with the Stokes beam (which is the only beam on axis) made separation of Stokes and pump easy. The highest input energy available in the Stokes beam was about 1 mJ in 0.3 cm^{-1} bandwidth with a 20 ns pulse duration. The best Stokes output energy achieved was 2 J which although representing a gain of 2000 (e^7) was insufficient to significantly deplete the pump. Under small signal conditions the Stokes power in the amplifier grows as $P_s = P_s(0) \exp(\gamma I_p L)$ and thus $\gamma I_p L \approx 7$ where γ is the Raman gain coefficient I_p the pump intensity and L the interaction length. The observed gain indicates that $\gamma \approx 0.2 \text{ cm GW}^{-1}$ at the cell operating pressure of 4 atmos CH_4 . This is considerably less than would be expected using the line centre value of $0.1 \text{ cm GW}^{-1} \text{ atmos}^{-1}$. It is necessary however to account for the effects of laser linewidth, pressure broadening of the Raman line and also pressure shift where the oscillator and amplifier are run at different pressures as was the case in this experiment. Measurements (C5.1) indicate the Raman width of the ν_1 transition of CH_4 varies with pressure as $\Delta\nu = 0.32 + 0.012 p \text{ cm}^{-1}$ with p in atmos. The line centre shifts with pressure at a rate $\nu - \nu_0 = 0.0175 p \text{ cm}^{-1}$. Assuming a Lorentzian gain profile, the peak gain for a shift of $\nu - \nu_0$ and a total Raman linewidth of $\Delta\nu = 0.32 + 0.012 p + \Delta\nu_L$ where $\Delta\nu_L$ is the laser linewidth is given by:

$$\gamma(\nu) = \frac{8 \times 10^{-3} p \Delta\nu}{(\nu - \nu_0)^2 + \left(\frac{\Delta\nu}{2}\right)^2} \text{ cm GW}^{-1}$$

Using this expression for a Raman oscillator pressure of 10 atmos and a Raman amplifier pressure of 4 atmos with a laser linewidth of 0.3 cm^{-1} we find a Raman gain of 0.17 cm GW^{-1} in reasonable agreement with experiment. Improvements in the oscillator performance (described in section C2) in terms of reduced linewidth and increased Stokes power indicate that significant pump depletion should be achievable under the same pump conditions.

C5.2 Model Calculations

During the past year a start has been made to extend the simple 1D model presented in the last annual report to include more realistic assumptions as

to the properties of the light guide. In the extended model a single pump beam moving at an angle to the Stokes beam and making reflections of finite reflectivity at the guide walls is considered. The model also considers the Stokes beam not to fill the light guide which is an important property if the Raman amplifier is to be used to amplify a number of beams travelling at different angles through the guide as occurs in the staged multiplexer. Typical results are shown in Figure C5.3 for the case where the area of the Stokes beam is 0.5 of the guide area. The curve for one reflection at $R = 85\%$ would be typical of the present experiment although larger fill factors than 0.5 can be obtained experimentally. With a low fill factor the effect of increasing the number of reflections is beneficial since the pump beam traverses the Stokes beam more often. The finite wall reflectivity can however keep the conversion efficiency down. For comparison the effect of multiple reflections from perfect reflectors is to greatly increase the conversion efficiency even when the fill factor is small.

C5.3 Materials for light guides

The present experiment uses float glass for the guide walls and obtains high reflectivity by operating at near grazing incidence ($< 1^\circ$). For a staged multiplexer however, operation over a larger acceptance angle will be necessary both as a consequence of the above discussion (since fill factors must be less than unity) but also from the point of needing to accommodate a large number of pump and Stokes beams. This rules out uncoated glass as a wall material due to its low reflectivity. The desired properties for the guide walls are flatness (so that pump beams are not focused), high reflectivity at grazing angles up to 10° or so and high damage threshold. In addition the coating should be inexpensive and applicable to large areas.

We have investigated some properties of the following materials: evaporated Aluminium, evaporated Rhodium, plated Rhodium, and silicon-coated float glass. All of these material possess low damage thresholds at normal incidence ($< 0.5 \text{ J cm}^{-2}$) however a number may be useful at near grazing incidence. Perhaps the most interesting material is silicon coated float glass (Pilkingtons spectraflo) since this is readily available in large

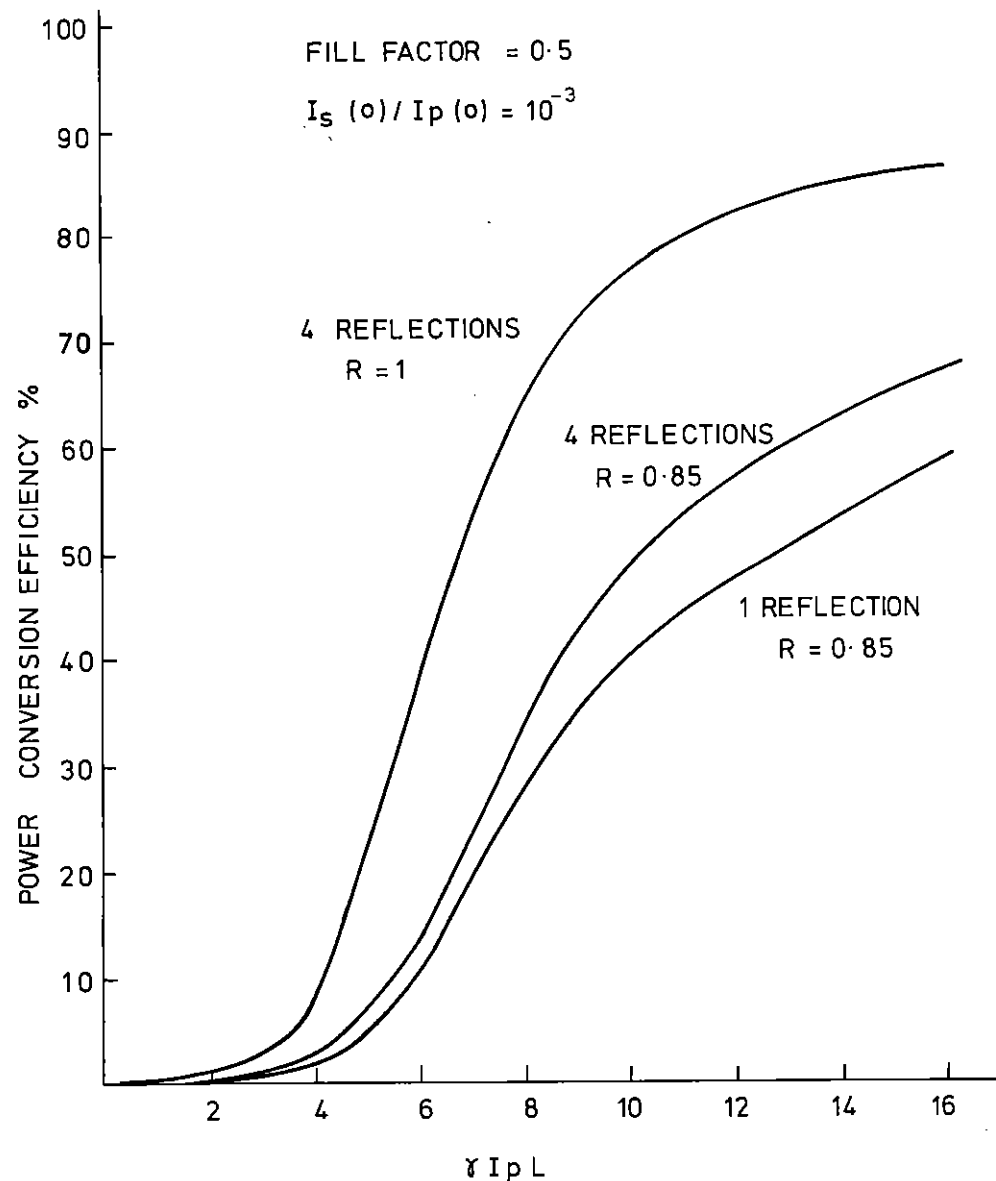


Fig. C5.3 Model calculations of light guided Raman amplifier efficiency.

sizes. The silicon coating is rather thin and since reflectivity at $0.25\ \mu\text{m}$ was not well known it was decided to measure the reflectivity using both a KrF laser and also the second harmonic of an Ar^+ laser at $256\ \text{nm}$.

The results are shown in Figure C5.4 which shows the reflectivity measurements as a function of angle. The solid curves are calculated from the Fresnel equations assuming $n = 4.4$ and $k = 0$. However the data are also fitted quite well with $n = 2.0$ and $k = 1.0$ but the data are fitted very poorly by $n = 1.7$, $k = 3.7$ the quoted values for single crystal silicon at $0.25\ \mu\text{m}$. Unfortunately for this thickness of silicon the reflectivity falls off rather too quickly with angle of incidence for spectrfloat to be of use as lightguide walls. The search continues for the most suitable material for this application.

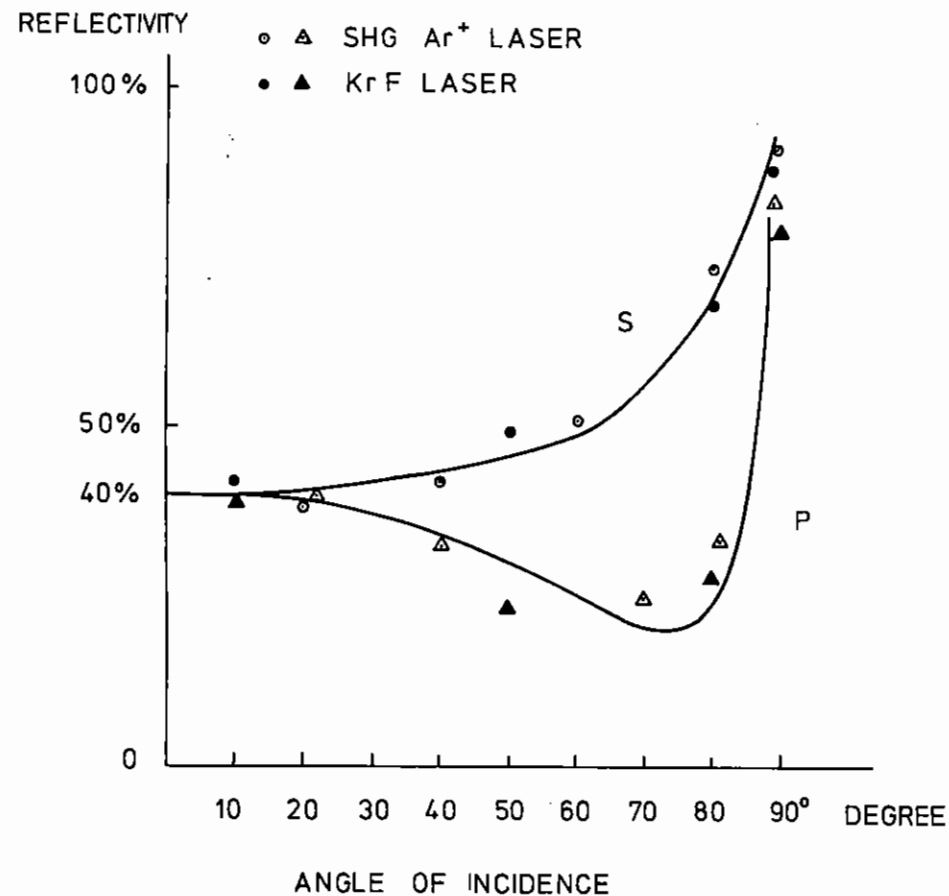


Fig. C5.4 Reflectivity of silicon-coated float glass at $0.25\ \mu\text{m}$.

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