

XUV AMPLIFICATION IN RECOMBINING
LASER PRODUCED PLASMAS

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ABSTRACT

In recent years (1985 – 1988) there has been a renewal of interest in X-ray and XUV (Extreme Ultra-Violet) lasers following the demonstration of amplification at 206 Å and 209 Å in a laser produced selenium XXIV plasma at the Lawrence Livermore National Laboratory. Recent progress in the field has been due to a combination of technical advances in large (Tera-Watt) pump lasers, optical engineering, fast time resolved XUV diagnostics, computer modelling and an increased understanding of the basic plasma and atomic physics processes involved.

This thesis reports on work carried out by the author on XUV laser schemes in the 200 – 50 Å spectral range. The work has included the construction of a novel time resolved XUV spectrometer, its absolute calibration using the Daresbury synchrotron source and its use in experiments on laser produced plasmas at the Rutherford Appleton Laboratory. Demonstrations of significant XUV gain at 182 Å in H like carbon plasmas, at 81 Å in H like fluorine plasmas and at 103, 105, 150 and 154 Å in Li like aluminium are reported. Comparisons of experimental observations with detailed numerical simulations are included.

Experiments on the optimisation of gain as a function of plasma energy and high Z radiative cooling for the H like fluorine scheme are described. The fluorine laser was a successful short wavelength scaling of the hydrogenic carbon scheme. However the scaling of the hydrogen like recombination scheme to still shorter wavelengths will require very high density plasmas. An experiment on the short pulse (3 ps) irradiation of aluminium foils at intensities up to 10^{17} W cm⁻² which resulted in fully ionised Al plasmas at temperatures of 400 eV and densities close to solid (10^{23} cm⁻³) is reported. Plasma conditions in this difficult to achieve high density regime will be needed before the hydrogenic scheme can be scaled into the water window (44 – 22 Å), an important spectral range yet to be reached.

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TO SYLVIA AND CLIFF
FOR MAKING ME CURIOUS !

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SYMBOLS AND ABBREVIATIONS

a_0	Bohr radius
A_{xy}	Einstein spontaneous emission coefficient
$A(m,l)$	Radiative decay rate
A_m	Atomic mass number
B_{xy}	Einstein stimulated emission coefficient
c	Speed of light
c_s	Ion sound speed
D	Illuminated length of line focus mirror
E	Absorbed plasma energy
E^*	Optimum plasma energy for gain
E_t	Optimum thermal energy for gain
E_i	Ionisation energy of plasma
E_{tot}	Total plasma energy $E_t + E_i$
E_H	Hydrogen like ion, ionisation energy
E_n	Energy of n th atomic level
E/L	Absorbed plasma energy per unit length
ΔE	Energy level separation $E_2 - E_1$
F	Focusing optics F number
$F(m,n,l)$	Absorption oscillator strength
g_n	n th level degeneracy
$g(\nu)$	Normalised line profile function
h	Planck's constant
$\hbar$	$h/2\pi$
$H_\alpha, \beta \dots$	Balmer series $2-n$ ($n \geq 3$) transitions
$He_\alpha, \beta \dots$	Helium like $1s^2-1snp$ ($n \geq 2$) transitions
I	Laser irradiance
I	Radiation flux
$I(\nu)$	Radiation flux as a function of frequency ν
J	Total angular momentum quantum number
K	Boltzmann's constant
$K(m,n,l)$	Collisional de-excitation rate
l	Orbital (angular momentum) quantum number
l	Plasma length
L	Line focus length
$Ly_\alpha, \beta \dots$	Lyman series $1s-np$ ($n \geq 2$) transitions

m	Magnetic quantum number
m	Lasant ion mass
M	Plasma mass per unit length
M^*	Optimum plasma mass for gain
M_b	Ablated plasma mass
n	Principle quantum number
n_c	Critical plasma density
n_e	Electron density
n_i	Ion density
n_i	Initial density of expanding plasma
n_f	Final density of expanding plasma
n_ℓ	Level nearest to thermal limit
n_s	Last level suppressed by continuum lowering
N_n	nth level electron population
ΔN	Population inversion density
ΔN_c	Critical population inversion density for > 0 gain
P	Electric dipole moment
$P_{\min}$	Minimum pump power for > 0 gain
r	Ratio of spontaneous to stimulated emission
r_{ab}	Stimulated absorption rate
r_{st}	Stimulated emission rate
r_{sp}	Spontaneous emission rate
r_0	Initial target radius
R	Plasma column radius
R	Line focus mirror radius
R_i	Expanding plasma initial radius
R_f	Expanding plasma final radius
S	Plasma source function ϵ/α
t	Target heating time $\approx$ laser pulse width
t_{photon}	Photon lifetime in a cavity
t_{spont}	Spontaneous emission lifetime $1/A$
T	Absolute temperature
T_i	Ion temperature
T_e	Electron temperature
T_i	Expanding plasma initial temperature
T_f	Expanding plasma final temperature
u_ν	Black body radiation field density

w	Line focus width
W'	Stimulated emission rate per electron
$W(\nu)$	Total transition rate as a function of frequency ν
Z	Atomic number
Z_i	Ionisation number of plasma
α	Gain coefficient
β	Input laser beam divergence
γ	Laser pulse shape factor
ϵ	Spontaneous emission coefficient
ζ	Adiabatic index
λ	Wavelength
ν	Frequency
ξ_0	Permittivity of free space
θ	Line focus mirror tilt angle
$\rho(\nu)$	Radiation field density as a function of frequency ν
τ	Radiation escape depth
χ_i	i th ionisation state, ionisation energy
ASE	Amplified Spontaneous Emission
APC	Active Pinhole Camera
KDP	Potassium Dihydrogen Phosphate $K_4H_2PO_4$
LLNL	Lawrence Livermore National Laboratory
LTE	Local Thermodynamic Equilibrium
NRL	Naval Research Laboratory
PET	Pentaerythritol $C(CH_2OH)_4$
RAL	Rutherford Appleton Laboratory
TAE	Target Area East
TIAP	Thallium Hydrogen Phthalate $TIHC_8H_4O_4$
TW	Tera-Watt
XUV	Extreme Ultra-Violet
ns	Nanosecond 10^{-9} sec
ps	Picosecond 10^{-12} sec
fs	Femtosecond 10^{-15} sec

SPECTROSCOPIC NOTATION

The majority of the spectral information in this thesis was taken from the Naval Research Laboratory report 7599 "*Atomic and Ionic Emission Lines below 2000 Å*" by R. L. Kelly and L. J. Palumbo, June 1973. Some of the spectroscopic notations and abbreviated forms used in this thesis are non-standard, but are widely used in the X-ray laser and plasma physics communities, and are described below. For spectra from plasmas containing various ionisation stages of a single element, the following nomenclature is used.

Ly_{α} , Ly_{β} etc. Lyman series lines of the form $1s-np$ ($n \geq 2$) of the hydrogenic ion. $Ly_{\alpha} = 1s-2p$, $Ly_{\beta} = 1s-3p$ and so on.

H_{α} , H_{β} etc. Balmer series lines $2-n$ ($n \geq 3$) of the hydrogenic ion. $H_{\alpha} = 2-3$, $H_{\beta} = 2-4$ and so on.

The fine structure of these lines is not generally included in the labelling of figures, as most of the fine structure components are well below the resolving power of the spectrometers used in our experiments. The effect of the fine structure on the gain in X-ray laser schemes is however included in some of the code simulations.

He_{α} , He_{β} etc. Helium-like ion transitions to the $1s^2$ ground state. $He_{\alpha} = 1s^2-1s2p$, $He_{\beta} = 1s^2-1s3p$ and so on.

For convenience when labelling lines from helium-like ions other than those to the $1s^2$ ground state, the $1s$ inner core is usually neglected, and the triplet and singlet states are denoted as follows :—

$He\ 3d^3-2p^3 \equiv 1s3d\ ^3D-1s2p\ ^3P$

$He\ 4d^1-2p^1 \equiv 1s4d\ ^1D-1s2p\ ^1P$ and so on.

Similarly for lithium like ions :—

$Li\ 3d^2-2p^2 \equiv 1s^23d\ ^2D-1s^22p\ ^2P$

$Li\ 3p^2-2s^2 \equiv 1s^23p\ ^2P-1s^22s\ ^2S$

Beryllium-like ions where there may be two p electrons are described by the more usual spectroscopic notation with the singlet and triplet states shown explicitly. Where multiple diffraction orders are observed, the order of the line is indicated by a number in front of the spectral description. In the case of plasmas containing more than one element, this is also followed by the symbol for the element, thus 3 Mg Ly_{α} refers to the 3rd diffracted order of hydrogen-like magnesium $1s-2p$.

CHAPTER 1. INTRODUCTION

The demonstration of the first red ruby laser (Light Amplification by Stimulated Emission of Radiation) by Maiman in 1960 [1] was closely followed by examples of laser action at other wavelengths in the visible spectral region. A very wide range of solids, liquids, gases and plasmas have since been made to lase using various methods of pumping. Commercial lasers are presently available at many wavelengths from the far infra-red, through the visible and into the ultra-violet.

Since the advent of the first optical laser work has been underway to develop lasers operating at ever shorter wavelengths. Figure 1.1 shows a brief history of the ability to generate coherent electric and electromagnetic waves. Many schemes for extreme ultra-violet (XUV), X-ray and even gamma ray lasers have been proposed. XUV and X-ray lasers are likely to have many important and unique applications in research and industry, but the search for such short wavelength lasers is hampered by many technical difficulties and basic physics problems not previously encountered in the development of optical lasers.

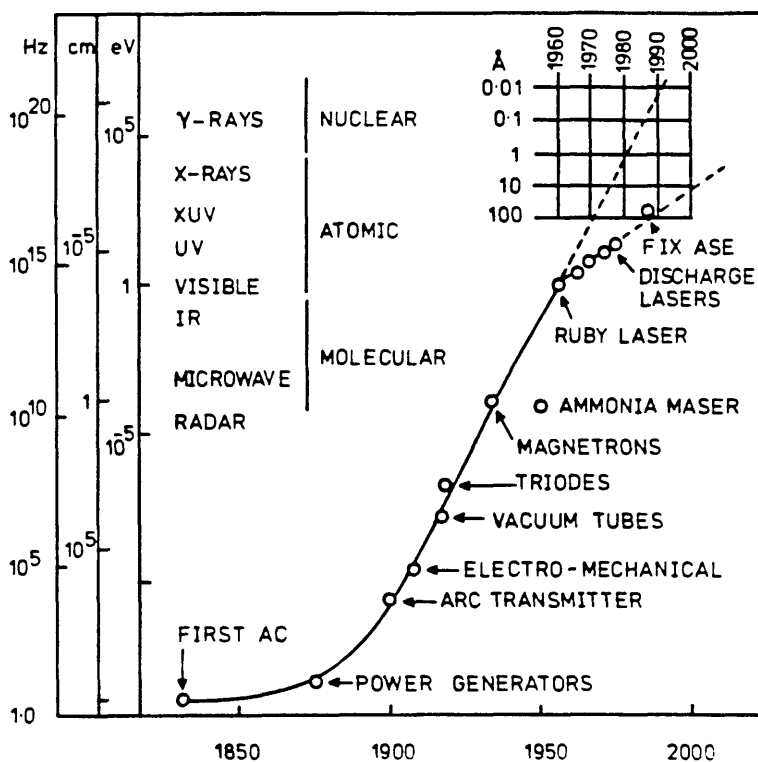


Figure 1.1 A history of the ability to generate coherent electric and electromagnetic waves.

Rapid progress has however been made in the field of XUV lasers during the last three years (1985 – 1988). The renewal of interest in X-ray and XUV lasers at this time followed the demonstration of gain at 206 Å and 209 Å in a laser produced Se^{24+} plasma [2,3] at the Lawrence Livermore National Laboratory.

1.1 WHY BUILD AN X-RAY LASER ?

The first optical laser was the result of pure research and was built with no particular application in mind. The unique and characteristic properties of spatial coherence, spectral purity, narrow line width, small beam divergence and high brightness have since led to the laser becoming an invaluable research tool in many branches of science. In application the laser has become an integral part of industries as diverse as ship building (precision measuring and laser welding), telecommunications (high bit rate optical fibre transmission) and semiconductor fabrication (laser annealing of silicon wafers). The laser has formed the basis of entirely new techniques such as holography, flash photochemistry and non-linear optics which have in turn opened up new avenues of research and new areas of physics.

As with the optical laser, it is impossible to envision all the eventual uses to which an X-ray laser could be put, but considering its unique output, likely high single shot brightness (10^{30} photons cm^{-2} sec^{-1} steradian $^{-1}$ in $\Delta\nu/\nu = 10^{-4}$) and sub-nanosecond pulse length a number of obvious examples spring to mind. The decrease in wavelength from the visible and UV to X-ray lessens diffraction effects and hence increases the resolving power of a laser used to probe a physical system. For example a fully coherent X-ray laser of a short enough wavelength might possibly be used to make holograms of individual molecules. For such an application the laser wavelength must be of the same order as the molecule being investigated; for a large, simple molecule, such as lead stearate this is ≈ 100 Å. The decrease of diffraction effects with wavelength would also make an X-ray laser a valuable tool for materials processing such as semiconductor fabrication, where structures are presently limited in size by the wavelength of UV light used to expose photoresists during manufacture.

The very high single shot brightness of an X-ray laser (as compared to sources such as synchrotrons, see figure 1.2) would allow short timescale events to be followed with an X-ray probe. An example of this is X-ray microscopy. At present biological samples must be exposed to synchrotron radiation for many minutes to record a transmitted X-ray image, and during this time the sample suffers radiation damage. A water window X-ray laser (44 – 22 Å where there is high contrast between carbon and oxygen in biological materials, (see figure 1.3) could be used to record X-ray images of living cells on a sub-nanosecond timescale. This is much shorter than that on which radiation damage could manifest itself. This technique could potentially be used to follow processes in living cells with ≤ 100 Å space and sub nanosecond temporal resolution. In this case the spatial resolution is limited by the electron microscope used to view the X-ray photoresist, rather than the wavelength of the laser.

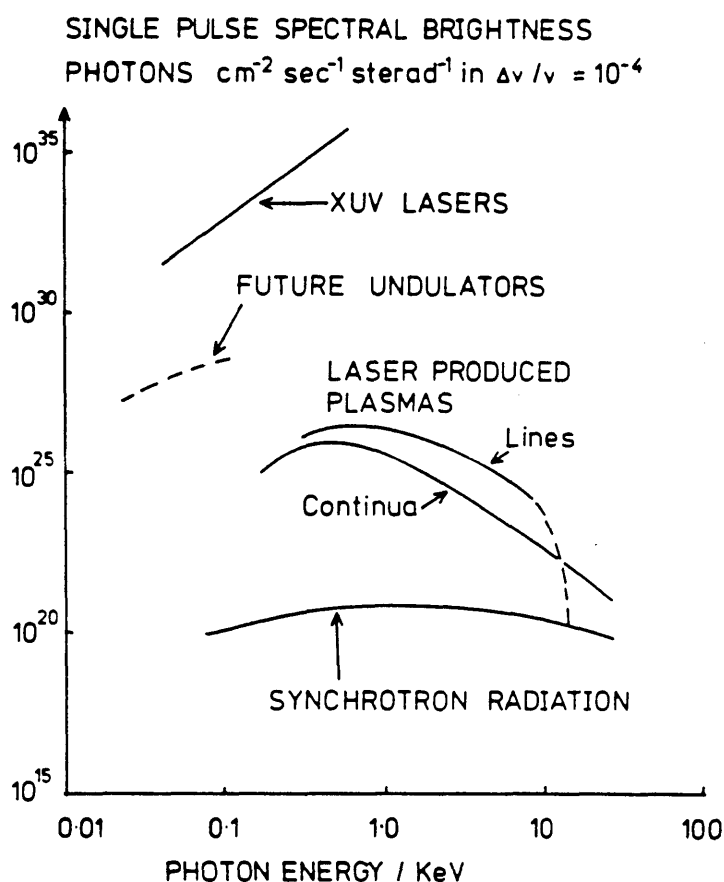


Figure 1.2 Single pulse X-ray brightness as a function of wavelength for several methods of X-ray production.

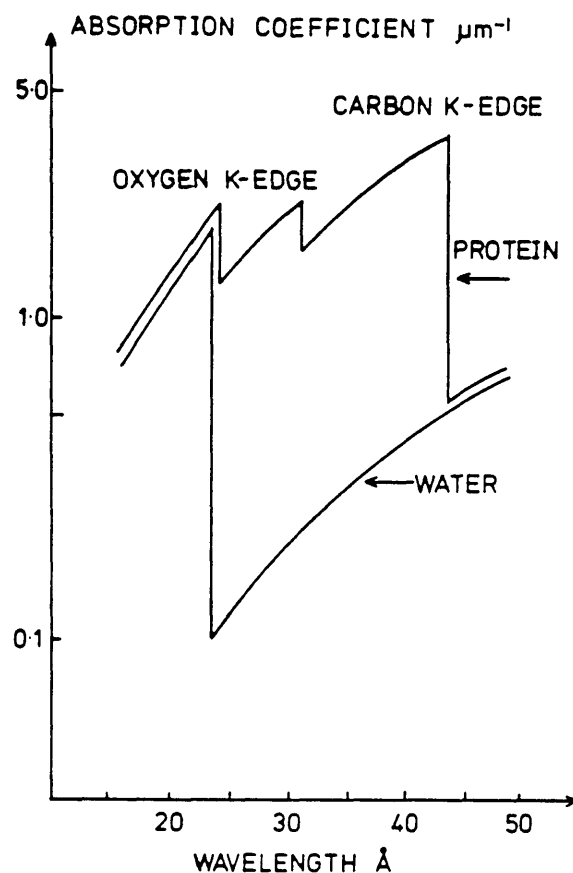


Figure 1.3 Differential XUV absorption by water and protein showing the "water window" between the carbon and oxygen K edges.

The high energy of X-ray photons and the narrow line width of an X-ray laser would have important applications in spectroscopy. Inner shell electrons could be directly excited, and this would provide a unique way of studying the atomic physics (excitation and relaxation processes) of many electron atoms on a sub-nanosecond timescale. At present, atoms have to be highly ionised to investigate most inner shell transitions, and this generates a large X-ray and XUV background as well as perturbing the atom.

X-ray laser research itself has also helped in the understanding of novel regimes of physics, and at the same time has provided the impetus for developments in fast time resolved X-ray and XUV diagnostics which have applications in other branches of science. The extremes of temperature and energy density required in X-ray laser gain media generate high fluxes of line and continuum X-ray emission. The study of X-ray emission from laser plasmas has provided valuable data for the improvement of the atomic physics models of many-electron atoms and ions.

The high pump laser powers ($> \text{TW}$) used in X-ray laser experiments lead to novel laser plasma interactions and instabilities. For instance a Bénard convective instability analogous to the atmospheric convection cells which drive much of our weather has recently been observed in a free standing thin foil laser target [4]. The investigation of such effects has resulted in an increased understanding of the interaction of powerful laser beams with matter, and generally stimulated the field of plasma physics.

1.2 PROBLEMS WITH X-RAY LASERS

Clearly to generate laser action in the X-ray region of the spectrum, a population inversion of a highly energetic electron transition must be used. Such transitions take place in the inner electron shells of atoms and ions. They are therefore difficult to excite and occur on timescales much shorter than transitions in the visible part of the spectrum. Note for example the ν^3 dependence of the rate r_{sp} for an electron transition given by [5] :—

$$r_{\text{sp}} = \frac{4\pi^3 \nu^3}{3\xi_0 hc^3} P^2$$

where ν is the frequency associated with the transition, P the electric dipole moment, ξ_0 the permittivity of free space, h Planck's constant and c the speed of light. The excitation of the required electron transitions demands very high temperatures and densities in the gain medium of an X-ray laser, typically greater than 10^5 K and 10^6 J g^{-1} . This in turn means that extremely large ($> \text{TW}$) pump powers are needed to run such a laser. It can be shown (see chapter 2) that the minimum pump power for a simple four state laser scales as ν^4 , and so pump power must increase dramatically when moving from optical to X-ray laser schemes.

Similarly, it can be shown (see chapter 2) that the minimum population inversion fraction for a laser scales as ν^2 or ν^3 depending on the line broadening mechanisms involved and again this causes severe problems when trying to operate an X-ray laser.

The high temperatures and consequent mass flow rates on the order of 10^7 cm s^{-1} also make it difficult to confine the gain medium and result in a correspondingly short plasma lifetime on the order of a few tens of nanoseconds. Optical lasers require relatively small gain length products ($\alpha \ell > 10^{-2}$ where α is the gain coefficient, ℓ the length of gain medium) in the amplifying medium to give a useful output [6], because it is easy to construct an optical cavity around the amplifier. A cavity effectively increases the length of the amplifier, and as the output is an exponential function $I = S(e^{\alpha \ell} - 1)$ of the path length ℓ through the amplifier this can have a dramatic effect on the output.

There are three major problems with building cavities for X-ray lasers. Firstly it is technically very difficult to construct a high efficiency normal incidence X-ray mirror, though in the last few years high / low Z multilayer structures which give back reflections by Bragg scattering off the layers have become available [7]. Secondly the inversion lifetimes of present day X-ray lasers are small (on the order of a few nanoseconds). Mirrors have to be placed tens of cm away from the gain medium to avoid being destroyed by reflected pump power before the inversion forms. This means that only a few passes would occur in an X-ray laser cavity ≈ 1 m long. Lastly, it is difficult to couple laser light out of an X-ray cavity. In contrast to the visible, X-ray mirror losses are due to absorption and scattering rather than transmission. This results in large amounts of energy being absorbed by the mirrors, which destroys them.

The difficulty in constructing an X-ray cavity has resulted in most X-ray laser experiments being based on amplified spontaneous emission (ASE). In such cases a long thin gain medium is used to give amplification along one preferred axis without the benefit of mirrors. However this limits the path length through the amplifier and consequently very high gains ($\alpha > 0.5 \text{ cm}^{-1}$, a factor of $10^2 - 10^3$ larger than for an optical laser) are needed to give measurable amplification from the few centimeter long plasmas typically produced in these experiments.

The large energy density requirements of an X-ray laser mean that high temperature, high density plasmas are the only suitable gain media presently available. These plasmas make inherently inefficient amplifiers as much of their energy resides in unwanted modes such as particle kinetic energy rather than in the inversion. Again there are quite severe problems in creating uniform, long scale length plasmas. The physics of such plasmas is not completely understood, and they are subject to a wide range of instabilities.

These instabilities result in poor pump energy coupling and produce density fluctuations which refract out of the plasma the X-rays that we wish to amplify. The atomic physics processes taking place in such plasmas (ionisation balance, opacity, radiation transport, etc) are also complex and difficult to model. Computer codes which incorporate the necessary plasma and atomic physics are being developed for use in simulating experiments and interpreting results. However these codes require data from experiments such as accurate line widths and wavelengths on which to base some of their calculations, and hence their development is an iterative process as both the codes and experiments progress.

There are also a number of purely technical problems associated with X-ray laser experiments. These include the difficulty in generating long, uniform plasmas and supplying sufficiently high pump powers to them. Present X-ray laser schemes only give amplification for a few nanosecond and fast time resolved diagnostics are needed to separate the ASE signal from the longer lived background plasma emission. The 200 – 50 Å spectral region in which current XUV and soft X-ray laser schemes operate is also difficult to work in, as these wavelengths are strongly absorbed by most materials.

1.3 RECENT TECHNICAL ADVANCES

The most successful X-ray laser experiments to date have been based on laser produced plasmas [2,8,9,10,11,35]. The large optical lasers recently developed for laser fusion research are capable of delivering the high pump powers required, and novel, high quality, line focus optics able to focus such lasers have become available. Good quality line foci in excess of 1 cm in length are needed to produce the long, thin, uniform plasmas around which to base a short wavelength ASE laser, and "traditional" cylindrical lenses have proved not to be ideal for this.

Lens aberrations from cylindrical lenses give rise to broad ($200\text{ }\mu\text{m}$) line foci, and techniques such as the use of spherical mirror and aspheric lens pairs have been developed to give the diffraction limited $25\text{ }\mu\text{m}$ wide line foci required for efficient heating of small X-ray laser targets.

Traditional XUV and X-ray reflection grating spectrometers disperse a spectrum onto an arc of the Rowland circle, and this makes them difficult to couple to time resolving elements. In recent years "flat field" XUV gratings have become available [12,13]. The line spacing is varied across these gratings in such a way that part of the Rowland circle is distorted into a straight line. This makes the coupling of XUV spectrometers to streak cameras relatively easy, and has lead to the construction of novel time resolved diagnostics capable of distinguishing an X-ray ASE signal from the long lived background plasma emission [14 and chapter 3]

There have also been major advances in computer modelling of the complex physics of laser produced plasmas and the atomic physics of high temperature plasmas containing multi-electron ions. Only recently have the effects of opacity, the kinetics of ionisation, pumping wavelength, energy absorption mechanisms and detailed hydrodynamics been appreciated and understood well enough to include in codes. Simpler codes gave plasma parameters widely different from experimental conditions which are critical for producing gain. These together with advances in basic plasma physics and atomic theory have made the planning and simulation of experiments more precise and resulted in a much better understanding of experimental results.

1.4 SOME POSSIBLE X-RAY LASER SCHEMES

A wide range of possible XUV and X-ray laser schemes have been suggested, and an equally wide range of pumping mechanisms proposed. This subject has been covered by a number of extensive reviews [15,16,17,18,19]. Laser produced plasmas are currently the most successful amplifiers below $200\text{ }\text{\AA}$. The laser plasma schemes split into three main areas, which are recombination, electron collisional excitation and photopumping. Recombination is discussed in more detail in chapter 2.

Briefly, in recombination schemes [8,9,10,35] a high temperature plasma cools and electrons recombine into highly excited states. The ion (H, Li and Na have been most successful) then de-excites through a radiative cascade. At some point in the cascade, a fast radiative transition depopulates a level (the lower laser level), generating a population inversion between this and the next highest level (upper laser level).

In collisional excitation [2,11], three-body collisions between ions (usually Ne or Ni like) and electrons in the plasma preferentially excite electrons into a level which then rapidly decays to the upper laser level, giving population inversion.

In photopump schemes, line or continuum emission from one plasma is used to pump a transition in a second plasma. The high temperature of the pump source should Doppler broaden the pump line, making the line matching conditions less critical than for the lower temperature pumps used in optical lasers. The advantage of the photopump scheme lies in the fact that a high effective excitation temperature can be achieved in the pumped plasma at relatively low temperatures and densities. This in turn reduces opacity and refraction losses in the amplifier.

To date a number of recombination and electron collisional excitation laser schemes have been operated in the $200 - 50 \text{ \AA}$ spectral region. Population inversion and gain [20,21,22] has also been observed in several photopump schemes, but at much longer wavelengths in the $1000 - 2000 \text{ \AA}$ regime.

Electron or ion beams are an alternative to laser beams for pumping X-ray lasers, but they are difficult to focus onto small targets and no working X-ray laser scheme powered in this way has been reported. Z pinches are an attractive method of generating long length, uniform, long lived (200 ns) plasmas which might be suitable for X-ray lasers [23], but no Z pinch based X-ray laser scheme showing gain has been reported.

Non linear optics has been used to successively frequency double and triple light from high power UV lasers down to $\approx 350 \text{ \AA}$, the seventh harmonic of a $\lambda = 2400 \text{ \AA}$ KrF excimer laser [24]. However this technique has not been successfully extended to the generation of XUV and X-ray light, as no suitable doubling medium is yet available for these wavelengths.

Free electron lasers can theoretically be tuned to any wavelength including X-ray, simply by altering the periodicity of the magnetic wiggler and the kinetic energy of the input electron beam. These devices are however still at an early stage of development.

Direct X-ray photopumping using a nuclear flashlamp has also been proposed, though this method is hardly applicable to laboratory X-ray laser experiments. While no detailed reports have appeared in open literature, it is believed that hard X-ray laser action at $\approx 14 \text{ \AA}$ has been observed in copper and iron rods pumped in this way [25,26].

It is interesting to note that while the excited states of inner shell electrons tend to be very short lived, some nuclei have excited states which are meta-stable for seconds or more. Such excited states again require nuclear reactions to pump them, but produce gamma rays when they decay. While a gamma ray laser or "graser" based on meta-stable excited nuclei is theoretically feasible, no evidence for gain on a gamma ray transition has been reported to date.

1.5 PRESENT STATE OF THE ART (Advances from 1985 to 1988)

The current interest in XUV and X-ray lasers is due in part to the recent demonstration of XUV gain in neon like Se^{24+} plasmas [2,3]. Foil targets consisting of $750 \text{ \AA}$ of Se evaporated onto a $1500 \text{ \AA}$ thick formvar substrate were irradiated in line focus geometry with two opposing beams from the Novette glass laser. Two 10 kJ , 450 ps pulses of green ($\lambda = 0.53 \text{ }\mu\text{m}$) light were focused to $2.2 \text{ cm} \times 200 \text{ }\mu\text{m}$ line foci using cylindrical lenses, to give a combined irradiance of $\approx 10^{14} \text{ W cm}^{-2}$ on target. The laser pulse heated the foil, with inverse Bremsstrahlung being the primary laser absorption mechanism. Neon like Se^{24+} ions were produced and monopole electron collisional excitation from the $2s^2 2p^6$ Ne like ground state to the $2p^5 3d$ and $2p^5 3p$ levels gave rise to a population inversion of the $2p^5 3s-2p^5 3p$ levels (see figure 2.4 for a simplified Ne like Se^{24+} level diagram).

NEON LIKE SELENIUM Se^{24+}

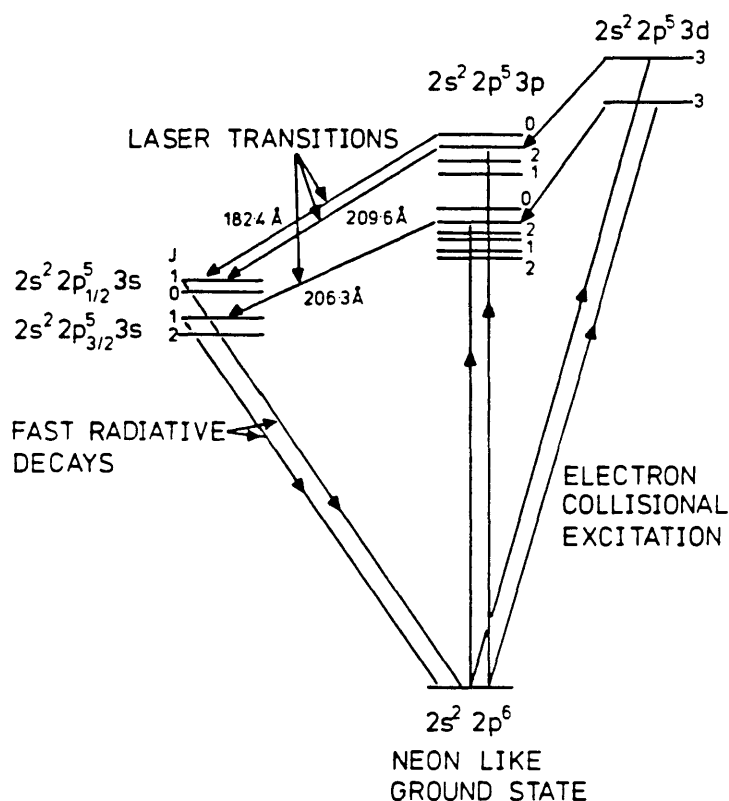


Figure 1.4 A simplified energy level diagram for the neon like selenium Se^{24+} X-ray laser scheme.

Time resolved transmission grating and grazing incidence spectrometers were used to record the X-ray emission spectra from the cylindrical plasma produced. Gains of $\approx 5.5 \pm 1.0 \text{ cm}^{-1}$ were seen on the two $J = 2-1$, $3p-3s$ transitions at 206.3 Å and 209.6 Å, but the $J = 0$ transition at 183 Å which was expected to show high gain was not initially observed. This transition has since been seen in longer length, Ne like Se^{24+} plasmas.

More recently the Ne like scheme has been scaled to shorter wavelengths by using higher Z lasing ions. Gains of 4 cm^{-1} have been seen for $3p-3s$ transitions at 155.0 Å and 157.1 Å in laser produced Ne like yttrium plasmas [27], and gains of 4 cm^{-1} seen at 106.4, 131.0, 132.7 and 139.4 Å in laser produced Ne like molybdenum plasmas [28]. In addition, Ne like copper and germanium amplifiers working at longer wavelength than the Se^{24+} laser have been investigated at the Naval Research Lab (NRL), to try to understand why the large gains predicted by codes are not seen on Ne like $J = 0$ transitions [29]. It has actually been suggested that recombination was an important pumping mechanism in the Se^{24+} laser [30], in an effort to explain the low gain on the $J = 0$ transitions.

Nickel like ions have a closed 3d shell analogous to the closed 2p shell of neon like ions (see figure 1.5). As with neon like ions, moderate gains have been observed for transitions in laser produced nickel like ions at the NRL. Again the population inversion mechanism is electron collisional excitation. In this case electrons are excited from the $3d^{10}$ ground state to the $3d4f$ and $3d4d$ levels, and laser transitions occur between the $3d4d$ and $3d4p$ levels. Small gains of $\approx 1 \text{ cm}^{-1}$ on $4p-4d$ transitions at wavelengths of 65.9, 71.0, 100.4 and 104.6 Å have been seen in Ni like europium plasmas produced by laser irradiation of thin foil targets [11].

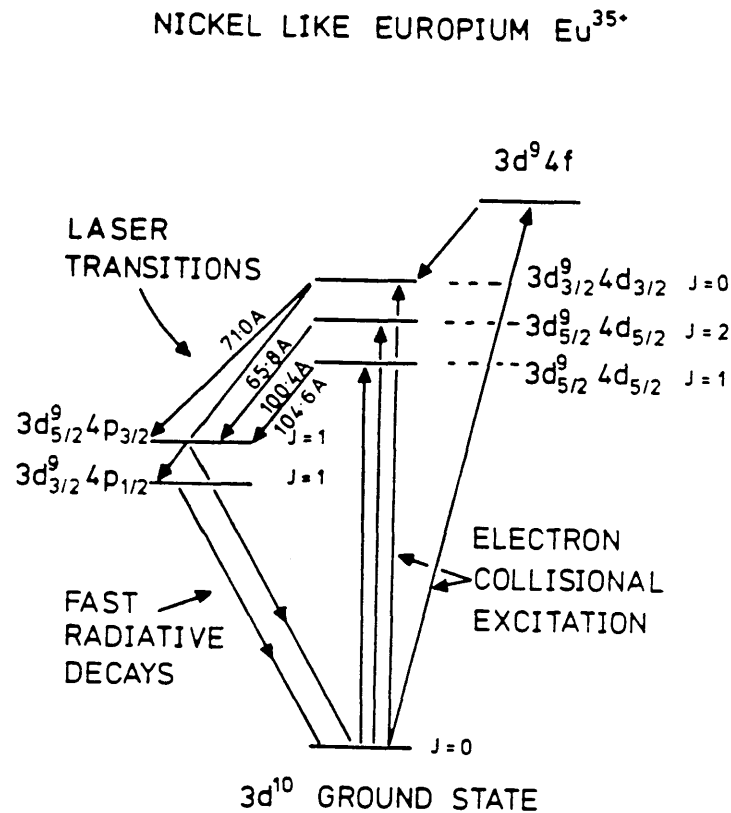


Figure 1.5 A simplified energy level diagram for the nickel like europium Eu^{32+} X-ray laser scheme.

A number of laser pumped recombination schemes have been highly successful in producing gain below 200 Å. These schemes have major energetic advantages over collisional excitation schemes (see figure 1.6) and consequently can be driven with a much smaller pump laser. For example a collisional excitation Se^{24+} laser operating at 206 Å requires a plasma temperature of $\approx 900 \text{ eV}$, compared to 10–20 eV for a C VI 182 Å recombination laser.

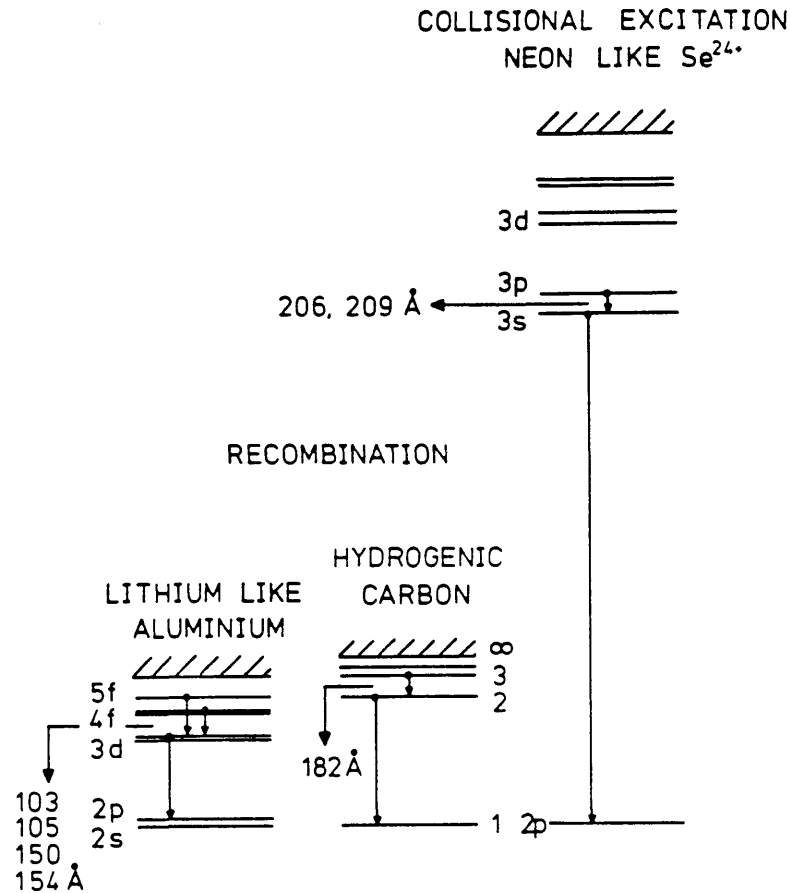


Figure 1.6 A comparison of the energy level spacing for some electron collisional and recombination X-ray laser schemes.

Figure 1.7 shows a simplified energy level diagram for a hydrogen-like carbon recombination laser. Population inversion of the $n = 3-2$ levels of laser produced C^{5+} ions at densities too low to support significant gain was seen as early as 1974 [31,32]. Small signal gains on the Balmer H_α transition at 182 Å were also reported by several groups [33,34]. However these experiments were carried out with very short (< 1 mm) plasmas and time integrated diagnostics, and as a result were not entirely conclusive. More recently gains of $4.0 \pm 1.0 \text{ cm}^{-1}$ in freely expanding carbon VI plasmas from thin fibre targets [8 and chapter 4] and gains of $5.5 \pm 1.0 \text{ cm}^{-1}$ in magnetically confined carbon VI plasmas [35,36] at 182 Å have been unambiguously demonstrated for longer (≈ 10 mm) lengths.

The thin fibre target recombination scheme has since been successfully scaled to shorter wavelengths by changing to a higher Z lasing ion. Lithium fluoride coated fibre targets were irradiated in line focus geometry and gain was seen at 81 Å in hydrogen-like fluorine [9 and chapter 5] using a pair of novel time resolved flat field XUV spectrometers [14 and chapter 3].

HYDROGENIC CARBON C^{5+}

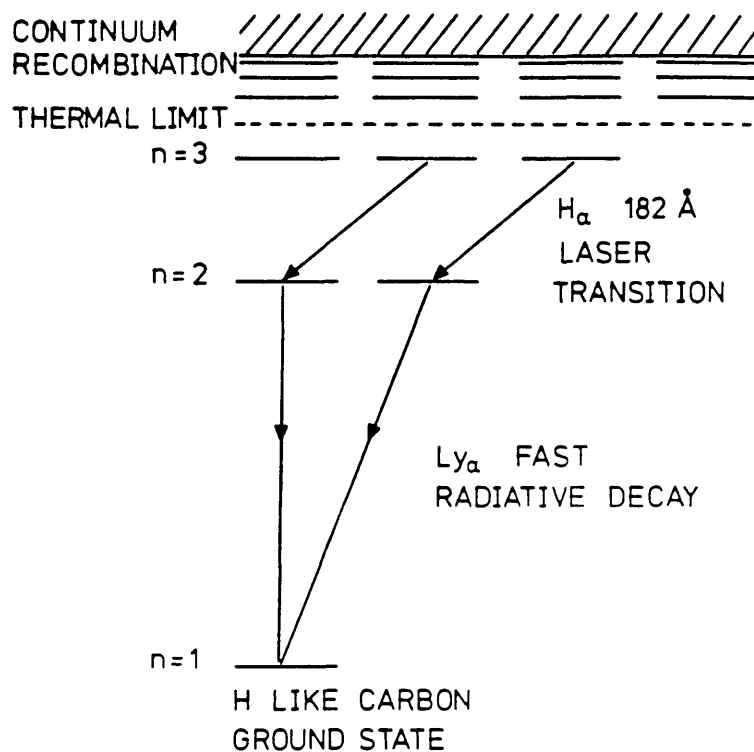


Figure 1.7 A simplified energy level diagram for the hydrogenic carbon recombination X-ray laser scheme

The effect of plasma energy variation on the fluorine IX amplifier was studied, and three gain regimes characterised by differing plasma energies and absorption ($-1.8 \pm 0.2 \text{ cm}^{-1}$), "high gain" ($4.4 \pm 1.0 \text{ cm}^{-1}$) and "low gain" ($2.0 \pm 1.0 \text{ cm}^{-1}$) were found [9 and chapter 5]. More recently, preliminary results from Osaka indicating low gain ($\approx 1-2 \text{ cm}^{-1}$) in H-like Na at $54 \text{ \AA}$ and H-like Mg at $46 \text{ \AA}$ from laser irradiated thin foils have been reported [37]. However these results were from time integrated diagnostics, and as such are not conclusive.

Lithium and sodium like ions have a single outer shell electron in the ground state, analogous to that for hydrogen like ions. They are expected to show gain on the $n = 4-3$, $5-3$ and $n = 5-4$, $6-4$ transitions respectively [38]. Gains in the region of $1-2 \text{ cm}^{-1}$ have been seen at 103.8 , 105.0 , 150.6 and $154.6 \text{ \AA}$ on the $5d-3p$, $5f-3d$, $4d-3p$ and $4f-3d$ transitions of lithium-like aluminium.

This effect has been seen in plasmas produced by line focus irradiation of both solid slab [39,40] and Al coated fibre targets [10 and Chapter 6]. The lithium-like recombination scheme has also been successfully scaled to shorter wavelength by changing the lasing ion from aluminium to sulphur. Experiments have also been carried out on a sodium-like copper recombination laser by the author but to date no evidence for gain has been found.

A joint recombination and photopump laser scheme has also been investigated recently by workers at RAL [41]. The $1s^2 2s^2 S_{1/2, 3/2} - 1s^2 5p^2 P_{1/2}$ transition in a lithium like calcium plasma from a laser irradiated fibre target was photopumped with $1s^2 2s^2 2p^3 P - 1s^2 2s 3d^3 D_3$ line radiation from a laser produced Be like manganese plasma. Code simulations showed no gain for the recombining calcium plasma alone, and small ($\approx 2.0 \text{ cm}^{-1}$) gains with the manganese flashlamp. Analysis of the data is presently underway, but there does not seem to be any evidence for significant gain.

There are a number of X-ray line and continuum overlaps between different ions, which might be suitable to base a pure photopump X-ray laser on [42]. Again some of these schemes have been investigated experimentally in laser produced plasmas (sodium neon scheme [43]) and also in Z pinch plasmas, but no gain in the soft X-ray region has been reported. In most cases this is thought to be due to insufficiently bright X-ray flashlamps. However, gains in the 1000 Å spectral range have recently been reported in Xe^{3+} and Kr^{3+} plasmas photopumped by soft X-rays [21,22]. A laser irradiated Ti rod produced the pumping radiation, which photo-ejected an inner electron from Xe^{2+} or Kr^{2+} ions. Auger decay from the $2+$ to $3+$ ionisation state resulted in a population inversion, and gains of $\approx 0.8 \text{ cm}^{-1}$ at 1089 Å in Xe and 907 Å in Kr were seen over lengths up to 9 cm.

1.6 DIRECTION OF PRESENT RESEARCH

There are now several X-ray laser schemes working in the 200 – 50 Å regime. All these schemes are based on laser produced plasmas and without exception they require very large energy inputs. There is now considerable interest in optimising these schemes, improving their efficiency to achieve saturation, and scaling them to shorter wavelength. A major goal still to be reached is a "water window" X-ray laser which would have important applications in biology and X-ray microscopy.

To increase the fraction of pump energy coupled into the target, various new target geometries are being tested. Recombination X-ray lasers based on thin ($\approx 7 \mu\text{m}$) fibres have been highly successful, but these targets are inherently inefficient and only absorb $\approx 10\%$ of the pump energy. Various thin foil targets have been used successfully in collisional excitation schemes, and similar targets are under consideration for recombination schemes. These targets have very good coupling efficiencies (80 – 90%), but foils supported along their full length suffer from problems with lateral heat transport removing energy from the narrow coated part of the target that we wish to heat. It is possible to make foil targets supported by thin cross pieces, but difficult to make such foils thin enough to explode and expand cylindrically, rather than simply ablating [44,45].

The interaction of matter with intense laser beams is not completely understood. A wide range of processes result in scattering of pump energy back out of the plasma, and generate instabilities which give rise to density fluctuations. This very complex area of physics is of considerable interest to workers in other fields, such as laser fusion and beat wave accelerators. A wide range of theoretical and experimental work is presently underway, to try to improve our understanding of this novel physical regime.

The theoretical optimisation of X-ray laser schemes in terms of target geometry, pump power, plasma mass, expansion rate and plasma energy has been considered in some detail, particularly for fibre based recombination lasers [46,47]. Experiments on the optimisation of such schemes in terms of plasma energy are now underway and have provided useful feedback to theoreticians for the improvement of codes [chapter 5].

The population inversion in recombination lasers is generated as the plasma is cooled by adiabatic expansion, radiation loss or interaction with a cold solid or gas. Cooling of the plasma is essential for the operation of the laser, and if the cooling could be enhanced then the population inversion might occur earlier in time and at higher density, resulting in a larger gain. The thin fibre recombination schemes under consideration use low to medium Z lasing ions and are cooled primarily by adiabatic expansion. Bremsstrahlung and radiative recombination loss from a plasma scales as Z^2 while line radiation loss scales as Z^4 , hence doping a low Z recombination laser with a small amount of high Z material should result in additional radiation cooling of the plasma.

A small enhancement of gain in hydrogenic carbon plasmas from a formvar foil coated in a thin layer of selenium has been reported [48]. In this experiment the adiabatic expansion of the plasma was the primary cooling mechanism and took place on a timescale of ≈ 500 ps. The selenium plasma had a cooling time of ≈ 100 ps and both enhanced the cooling of the carbon plasma and increased electron density which enhanced the recombination without increasing the C^{5+} ion density and hence opacity. An experiment on radiative cooling of hydrogenic fluorine plasmas by a range of higher Z dopants has been carried out by the author and is described in chapter 7.

The cooling of a recombining plasma may also be enhanced by interaction with a solid, and this effect has been reported for magnetically confined, radiation cooled hydrogenic carbon plasmas produced by long pulse (75 ns) irradiation of disc targets with a CO_2 laser [35]. A factor of 4 enhancement in axial $182 \text{ \AA}$ emission was found when an array of carbon fins was placed at the edge of the plasma column.

Multi-layer X-ray mirrors have recently become available. These devices are composed of layers of high and low Z material such as carbon and platinum deposited on a silicon wafer. As with dielectric mirrors for the optical, they are built to reflect one particular wavelength. Single mirrors and cavities have been used to give double and triple pass amplification in collisionally excited, laser produced Se^{24+} plasmas and resulted in a factor of 2 enhancement of the 206 and $209 \text{ \AA}$ laser lines [49]. A major problem is how to couple the laser light out of the cavity. In the optical, loss from a mirror is generally due to transmission and so one mirror of a laser cavity can be made partially transmitting to allow out some of the laser light.

Conversely for X-ray mirrors the principal loss mechanisms are absorption and scattering, and to couple out laser light by transmission a very thin (few 10^3 atoms) thick mirror with no substrate would be required. An alternative method would be to use a multilayer mirror with a stepped surface. The steps would act as a reflection diffraction grating, with the zero order feeding back into the cavity and laser output via the first and higher order diffracted beams.

Neon and nickel like electron collisional schemes and hydrogen and lithium like recombination schemes have all been scaled to shorter wavelength by increasing the atomic number of the lasing ion. However this process cannot be continued indefinitely to ever shorter wavelengths. The energy requirements of these schemes increases rapidly with Z and existing X-ray lasers already operate at the limit of the available pump power sources. There are also a number of physical as well as technical limits to the short wavelength scalability of these schemes. For instance the initial density required to achieve maximum gain in the fibre based H like recombination X-ray laser scales as Z^7 [chapter 2]. Obviously the initial density required rapidly becomes much greater than solid, and as a result this scheme is unlikely to operate beyond $Z = 13$ (aluminium) at a wavelength of 39 Å.

Due to this sort of limitation, new schemes, such as the sodium like recombination laser, which are theoretically more suitable for shorter wavelength amplifiers are presently being investigated. Recent technical advances in chirping and other pulse compression techniques have resulted in optical and UV lasers operating at ultra short (femtosecond) pulse lengths. These may be ideal pump sources for shorter wavelength X-ray lasers as they can produce plasmas at close to solid density on very short timescales [50]. The interaction of such short pulse lasers with matter is a completely new area of physics and is currently under investigation by both theoreticians and experimentalists. Chapter 8 describes an experiment on the short pulse (3 ps) ionisation of solid aluminium targets using the SPRITE KrF laser system and reports observations of ≈ 350 eV hydrogen like Al plasmas at near solid densities.

Short pulse pump lasers also appear to be good candidates for driving photopump X-ray lasers at short wavelengths, as they generate high fluxes of X-rays when focused onto solid targets. They may also be used to pump more exotic schemes such as traveling wave excitation where a wedge shaped lens is used to heat a fibre or foil target in a zone which moves along the target in step with the ASE pulse [51], thus lowering the pump energy requirements

Finally, applications of existing X-ray laser schemes are being investigated. The 200 – 50 Å spectral region is not particularly useful to work with as it is heavily absorbed by most materials. Despite this, work towards achieving spatially coherent output from a laser plasma based ASE amplifier is already underway [52], and some preliminary experiments on imaging of wires placed in the output of a Se^{24+} laser have recently been carried out [53].

1.7 THE AUTHOR'S CONTRIBUTION TO THIS WORK

The experiments described in this thesis were on a scale which required the active participation of several experimentalists and theoreticians. The author made significant contributions to either the planning, diagnosis or analysis of all the work reported in this thesis. A detailed list of co-workers and their areas of interest is included in appendix I.

The author contributed to the design and was responsible for the construction and initial run up of the "transverse" time resolved flat field XUV spectrometer used on X-ray laser experiments at RAL. The author played a major experimental role in the spectrometer calibration at the Daresbury SRS [chapter 3], and set up and ran at least one primary diagnostic on the axial / transverse hydrogenic carbon X-ray laser [chapter 4], fluorine X-ray laser [chapter 5], lithium like aluminium [chapter 6], radiative cooling [chapter 7] and picosecond spectroscopy [chapter 9] experiments. He was the main investigator in the radiative cooling and axial / transverse carbon work and made significant contributions to the analysis of all the above experiments. The author was the major contributor to the axial / transverse carbon analysis, fluorine X-ray laser analysis and radiative cooling analysis and modeling. This included the writing of the paper on the optimisation and modelling of the fluorine laser submitted to Optics Communications, and the presentation of some of the above work at various international conferences.

The majority of the experiments reported in this thesis represent significant contributions to the fields of X-ray lasers or laser plasma interactions. Most experiments were collaborations between various British universities and research laboratories. A number also involved international collaborations with groups from France, Japan and the USA. The work in this thesis is supported by a number of refereed articles [8,9,10,50], unrefereed reports for the Rutherford Appleton Laboratory [112 to 116] and published conference proceedings [18,19,91, and 117 to 126]. A number of additional publications based on work reported in this thesis are presently in preparation.

CHAPTER 2. THEORY

This section will concentrate on the theory of fibre based hydrogenic recombination XUV lasers. We will examine their mode of operation, optimisation in terms of plasma parameters, pump power, pump wavelength and plasma conditions, and their scaling to shorter wavelengths. Before considering these points in detail, it is instructive to review some general laser theory. This can give some useful insights into the physical processes involved in lasers and the inherent difficulty of constructing an X-ray laser.

2.1 GENERAL LASER THEORY AND SHORT WAVELENGTH SCALING.

Let us consider a simple two state system with energy levels E_1 and E_2 , degeneracy g_1 and g_2 and electron populations N_1 and N_2 . The rates of spontaneous and stimulated single electron transitions between the lower level E_1 and upper level E_2 are given by the Einstein A and B coefficients [5].

$$\text{Rate of spontaneous emission } r_{sp} = N_2 A_{21} \quad (2.1)$$

$$\text{Rate of stimulated emission } r_{st} = N_2 B_{21} \rho(\nu) \quad (2.2)$$

$$\text{Rate of stimulated absorption } r_{ab} = N_1 B_{12} \rho(\nu) \quad (2.3)$$

where $\rho(\nu)$ is the radiation field density as a function of frequency ν .

To find an expression for the radiation field, we use the conservation of energy and the Boltzmann equation. From the conservation of energy it follows that $r_{sp} + r_{st} = r_{ab}$ in a closed system. Substituting (2.1), (2.2) and (2.3) into this expression gives :-

$$\rho(\nu) = \frac{N_2 A_{21}}{N_1 B_{12} - N_2 B_{21}} \quad (2.4)$$

For a system in thermal equilibrium, the ratio of the populations of the two levels is given by the Boltzmann distribution [5].

$$N_2/N_1 = g_2/g_1 \exp(-h\nu/KT) \quad (2.5)$$

where h is Planck's constant, K Boltzmann's constant, T the absolute temperature, and the photon energy $h\nu = \Delta E$.

Setting $g_1 = g_2$ for simplicity and substituting (2.5) into (2.4) we find :-

$$\rho(\nu) = \frac{A_{21}}{B_{12} \exp(h\nu/KT) - B_{21}} \quad (2.6)$$

In fact B_{21} the probability of stimulated emission from E_2 to E_1 is equal to B_{12} the probability of a photon being absorbed by an electron in the E_1 level. Absorption is the inverse process to stimulated emission, thus we may set $B_{12} = g_1/g_2 B_{21}$ and the radiation field becomes :-

$$\rho(\nu) = \frac{A_{21}}{[\exp(h\nu/KT) - 1]B_{21}} \quad (2.7)$$

By definition, this must be identical to the black body radiation field u_ν

$$u_\nu = \frac{8\pi h \nu^3}{c^3} \frac{1}{\exp(h\nu/KT) - 1}$$

$$\text{Thus } \frac{A_{21}}{B_{21}} = \frac{8\pi h \nu^3}{c^3} \quad (2.8)$$

It is important to note the ν^3 dependence of the ratio of the Einstein spontaneous to stimulated emission coefficients.

It is now simple and instructive to calculate the ratio of spontaneous to stimulated emission r for a black body.

$$\begin{aligned} r &= r_{sp} / r_{st} \\ &= \frac{N_2 A_{21}}{N_2 B_{21} \rho(\nu)} = \exp(h\nu/KT) - 1 \end{aligned} \quad (2.9)$$

Using this expression for r at $T \approx 300$ Kelvin we see that for microwave transitions with a wavelength $\lambda \approx 0.3$ m, $r \approx 1.5 \times 10^{-4}$ and that most of the radiation emitted by such sources is stimulated and coherent. For green light at $\lambda \approx 6000$ Å, $r \approx 10^{35}$ and there is practically no stimulated emission, while for soft X-rays with $\lambda = 100$ Å, r is on the order of 10^{135} !

However equation (2.8) assumes that the system is in thermal equilibrium and $N_2 < N_1$ as required by the Boltzmann equation. If we can destroy the equilibrium and achieve a population inversion where $N_2 > N_1$ then stimulated emission in the visible or at even shorter wavelengths becomes possible. This can be very difficult to do, and we shall now examine the threshold conditions and frequency dependence for the formation of a population inversion and pump power required to achieve gain in an inverted system.

2.2 THE THRESHOLD POPULATION INVERSION.

It is useful to consider the possible loss mechanisms and gain in a simple laser and relate this to the minimum population inversion which must be achieved before these losses are overcome. Our laser will consist of an amplifying medium placed between two mirrors to form a cavity. The most important loss mechanisms are :—

- a) Opacity, absorption in the amplifier due to energy levels not directly involved in the laser transition, as no real atomic system has just two levels.
- b) Scattering by optical defects in the gain medium, such as high density gradients in a plasma.
- c) Diffraction losses by apertures, including the ends of the gain medium.
- d) Transmission, absorption and scattering by the mirrors used to make a cavity and the actual output of the laser.

These losses can be described by a single parameter t_{photon} , the lifetime of a photon within the cavity. For a mirrorless soft X-ray laser based on single pass ASE in a laser produced plasma there will be no feedback into the gain medium and the upper limit of t_{photon} will be set by the time for a photon to travel the full length ℓ of the gain medium. Thus $t_{\text{photon}} \leq \ell/c$ which is on the order of 3×10^{-11} for a 10 mm plasma. For our simple two level system we write the stimulated transition rate per electron in E_2 as :—

$$W' = \rho(\nu) B \quad (2.10)$$

Substituting in the expression for B from (2.8) gives :—

$$\frac{W'}{A_{21}} = \frac{\rho(\nu) c^3}{8\pi h \nu^3} \quad (2.11)$$

In fact this expression is inaccurate as real atomic systems will always emit lines with finite widths. For laser lines the line width is generally narrower than the Doppler width, and may approach the natural atomic transition line width [6]. We can take account of this with the line profile function $g(\nu)$, which is normalised such that $\int_0^\infty g(\nu) d\nu = 1$. $g(\nu) d\nu$ is the probability that a photon between ν and $(\nu+d\nu)$ will stimulate an emission. Equation (2.11) becomes

$$W' g(\nu) d\nu = \frac{\rho(\nu) c^3}{8\pi h \nu^3} \cdot A_{21} g(\nu) d\nu \quad (2.12)$$

To model radiation transfer through the system we now introduce the radiation flux or intensity $I(\nu)$ given by $I(\nu) = c\rho(\nu)$, and the spontaneous emission lifetime $t_{\text{spont}} = 1/A_{21}$. Integrating both sides of (2.12) we obtain the total transition rate $W(\nu)$ as a function of frequency due to a monochromatic beam of radiation, frequency I_ν passing through the gain medium.

$$W(\nu) = \frac{c^2}{8\pi h \nu^3 t_{\text{spont}}} \cdot g(\nu) I_\nu \quad (2.13)$$

The net rate of change of energy per unit distance the beam travels, due to stimulated transitions both up and down is $dE/dx = h\nu (N_2 - N_1 g_2/g_1) W(\nu)$, where we have taken account of the degeneracy of E_1 and E_2 . The rate of increase of intensity of the beam in a small time interval is then :-

$$\frac{dI(\nu)}{dt_{\text{gain}}} = c h \nu (N_2 - N_1 g_2/g_1) W(\nu) \quad (2.14)$$

For a population inversion, $N_2 \geq N_1 g_2/g_1$ thus (dI/dt) is positive and we have gain, an increase in intensity. It is convenient to write $(N_2 - N_1 g_1/g_2) = \Delta N$ the population inversion density. If $\Delta N \leq 0$ then the laser medium will simply absorb the radiation that we are trying to amplify. For simplicity we define the total loss rate using t_{photon} :-

$$\frac{dI(\nu)}{dt_{\text{loss}}} = \frac{I_\nu}{t_{\text{photon}}} \quad (2.15)$$

Clearly for there to be a net gain :-

$$\frac{dI(\nu)}{dt_{\text{gain}}} - \frac{dI(\nu)}{dt_{\text{loss}}} \geq 0 \quad (2.16)$$

Substituting the expressions for dI/dt gain and loss (2.14) and (2.15) into (2.16) gives a minimum requirement for the population inversion density ΔN needed for there to be gain on a line centered on ν_0 with a lineshape $g(\nu_0)$:-

$$\Delta N = (N_2 - N_1 g_2/g_1) \geq \frac{8\pi t_{\text{spont}} \nu_0^2}{c^3 g(\nu_0) t_{\text{photon}}} \quad (2.17)$$

We have to know something about $g(\nu_0)$ in order to use the inequality (2.17). The lineshape will be determined by the broadening mechanisms in the gain medium [6]. In an X-ray laser based on a high temperature laser plasma, Doppler broadening will be the dominant mechanism and will give rise to a Gaussian line profile [54]. This profile may be slightly modified, for instance in lower Z materials there will be a contribution from Stark broadening which will effect the outer wings of the profile, while fine structure splitting may effect the symmetry of the profile. These effects can become important when carrying out detailed numerical simulations, where a Voigt profile (convolution of Stark and Doppler profiles) is more appropriate. However for our simplistic argument a Gaussian is a good approximation to the line shape [55].

$$g(\nu) = \frac{2(\pi \ln 2)^{\frac{1}{2}}}{\pi \Delta \nu} \cdot \exp \left[-\frac{1}{2} \left[\frac{\nu - \Delta \nu}{\Delta \nu} \right]^2 \right] \quad (2.18)$$

For a line centered on ν_0 :-

$$g(\nu) = \frac{2(\pi \ln 2)^{\frac{1}{2}}}{\pi \Delta \nu} \quad (2.19)$$

$$\Delta \nu = \frac{2\nu_0}{c} \left[\frac{2KT \ln 2}{m} \right]^{\frac{1}{2}} \Rightarrow g(\nu) = \frac{c}{\pi \nu_0} \left[\frac{m \pi}{2 K T} \right]^{\frac{1}{2}} \quad (2.20)$$

Where m is the mass of the lasant ion. Substituting (2.20) into (2.17) we find an expression for the minimum population inversion density :-

$$(N_2 - N_1 g_2/g_1) \geq \nu_0^3 \cdot \frac{8\pi^2}{c^4} \left[\frac{t_{\text{spont}}}{t_{\text{photon}}} \right] \left[\frac{2KT}{\pi m} \right]^{\frac{1}{2}} = \Delta N_c \quad (2.21)$$

From (2.21) it is immediately obvious that building an X-ray laser will be difficult, as the minimum population inversion needed for there to be any gain scales as ν^3 . In addition, for a mirrorless single pass ASE amplifier t_{photon} will be small and T must be large to excite the necessary X-ray transitions, further increasing ΔN_c .

2.3 MINIMUM PUMP POWER.

Instead of a simple two level system, we now consider a more realistic four-level laser with ground state $n = 1$ and a laser transition between the $n = 3$ and $n = 2$ levels. Starting with the $n = 2, 3$ and 4 levels essentially empty, only a few electrons need to be pumped into the $n = 3$ level to give a $3-2$ inversion. For the threshold for gain, it follows that N_3 must be at least ΔN_c . The energy loss from the $n = 3$ level by spontaneous emission is $N_3 h\nu A_{32}$ and this must be exceeded by the minimum pump power P_{min} .

$$P_{\text{min}} = \frac{\Delta N_c h\nu}{t_{\text{spont}}} \quad (2.22)$$

Substituting our expression (2.21) for ΔN_c into (2.22) gives :-

$$P_{\text{min}} = \frac{\nu_0^4 8\pi^2 h}{c^4 t_{\text{photon}}} \left[\frac{2KT}{m\pi} \right]^{\frac{1}{2}} \quad (2.23)$$

In a realistic experiment we might imagine a $100 \text{ \AA}$ ASE "laser" consisting of a single passed 10 mm long low Z ($Z = 9$) laser produced plasma. The upper limit of t_{photon} will be set by the transit time through the gain medium ($\ell/c \approx 3 \times 10^{-11} \text{ s}$) and KT will typically be a few tens to hundreds of eV for a laser plasma. Equation (2.23) gives a value of $P_{\text{min}} = 3 \times 10^9 \text{ W cm}^{-3}$ for the volumetric pump power. Remember that this is for the break even case where $\Delta N = \Delta N_c$, giving $\alpha = 0$. Taking into account that much of the pump energy will go into modes other than the $n \rightarrow 3$ transitions which populate the upper laser level, we can expect to increase this estimate by several orders of magnitude for there to be measurable gain. With an optimistic 10% pump laser coupling efficiency into the target this easily brings the experiment into the Tera-Watt pump power regime.

2.4 RADIATION TRANSPORT AND GAIN

To find an expression for the gain of our simple system, we again consider a beam of monochromatic radiation $I(\nu)$ passing through a volume element of unit cross section and length dx of the gain medium. We note that dt is simply dx/c and so the energy loss through absorption is given by :-

$$\frac{dI(\nu)}{dx_{\text{loss}}} = N_1 B_{12} I(\nu) g(\nu) h\nu/c \quad (2.24)$$

Similarly the energy emitted is the sum of the spontaneous and stimulated emissions

$$\frac{dI(\nu)}{dx_{\text{gain}}} = N_2 A_{21} g(\nu) h\nu + N_2 B_{21} I(\nu) g(\nu) h\nu/c \quad (2.25)$$

summing (2.24) and (2.25) gives an expression for the radiation transfer equation

$$\begin{aligned} \frac{dI(\nu)}{dx_{\text{tot}}} &= \frac{dI(\nu)}{dx_{\text{gain}}} - \frac{dI(\nu)}{dx_{\text{loss}}} \\ \frac{dI(\nu)}{dx} &= h\nu g(\nu) \left[N_2 A_{21} + \frac{I(\nu)}{c} \left[N_2 B_{21} - N_1 B_{12} \right] \right] \end{aligned} \quad (2.26)$$

Remembering that $B_{12} = g_2/g_1 B_{21}$ and using (2.8) to give B_{21} in terms of A_{21} , the radiation transfer equation (2.26) becomes :-

$$\frac{dI(\nu)}{dx} = h\nu A_{21} \left[N_2 + \frac{I(\nu) c^2}{8\pi h\nu^3} \left[N_2 - g_2/g_1 N_1 \right] \right] \quad (2.27)$$

we write the volume emission coefficient $\epsilon(\nu)$ which gives the spontaneous energy per unit volume in unit time as :-

$$\epsilon(\nu) = h\nu A_{21} N_2 g(\nu) \quad (2.28)$$

and similarly we write the volume gain coefficient $\alpha(v)$ as :-

$$\alpha(v) = \frac{\lambda^2}{8\pi} g(v) A_{21} \Delta N \quad (2.29)$$

where ΔN is the population inversion density ($N_2 - g_2/g_1 N_1$). The radiation transfer equation can then be written more simply as :-

$$dI(v)/dx = \epsilon(v) + \alpha(v) I(v) \quad (2.30)$$

If we assume that $\epsilon(v)$ and $\alpha(v)$ are independent (the gain does not saturate and significantly effect the radiation field), and that the gain medium is uniform, then (2.30) can be integrated to give an expression for the intensity of a beam after passing through a length ℓ of the gain medium.

$$I(v, \ell) = I(v, 0) e^{\alpha(v)\ell} + \frac{\epsilon(v)}{\alpha(v)} \left[e^{\alpha(v)\ell} - 1 \right] \quad (2.31)$$

If $I(v) = 0$ at $x = 0$ as in a single pass ASE laser then (2.31) becomes :-

$$I(v, \ell) = \frac{\epsilon(v)}{\alpha(v)} \left[e^{\alpha(v)\ell} - 1 \right] \quad (2.32)$$

It is usual to write $\epsilon(v)/\alpha(v) = S(v)$ the source function. For a single pass ASE gain medium of length ℓ we expect the output to be :-

$$I(v, \ell) = S(v) \left[e^{\alpha(v)\ell} - 1 \right] \quad (2.33)$$

Clearly for an optically thin plasma with no gain $\alpha(v) \rightarrow 0$ and it is easy to see that (2.33) reduces to $I(v, \ell) = \epsilon(v)\ell$ which is simply the spontaneous emission signal from a plasma length ℓ of unit cross section. If $\alpha(v)$ is positive then $I(v, \ell)$ increases exponentially with length and the laser has gain, while if $\alpha(v)$ is negative the plasma will become optically thick and absorb radiation. From (2.29) we can also see that for a given population inversion density ΔN , the gain scales as λ^2 and so at shorter wavelengths the gain decreases unless the population inversion density is increased accordingly.

2.5 RECOMBINATION LASERS.

The recombination scheme was first proposed by Gudzenko and Shelepin in 1965 for use in the visible spectral region [56]. Recombination XUV and X-ray lasers have been the subject of theoretical studies for some time, and have been particularly successful in producing gain below 200 Å in the last three years (1985 – 1988). The principle of the recombination laser is rather simple. A plasma produced by laser irradiation of a target, gas discharge, electron beam etc is allowed to cool by adiabatic expansion, radiation loss or interaction with a cold heat sink. As the plasma cools, electrons recombine via three body processes into high lying electronic levels. These excited states then decay via collisional or radiative transitions to lower levels, and usually cascade through a series of intermediate states. For closely spaced states, collisional transitions dominate, while larger energy transitions decay predominantly by radiative mechanisms. This follows from the collisional and radiative transition rates. Decays through both these channels are possible and the recombination cascade is thus described by the collisional–radiative model [57]. The de–excitation rate coefficients for collisional processes [58] in $\text{cm}^3 \text{s}^{-1}$ can be written as :–

$$K(m, n) \approx \frac{F_{mn} \exp \left[\frac{-\Delta E_{mn}}{K T_e} \right]}{\Delta E_{mn} T_e^{\frac{1}{2}}} \quad (2.34)$$

$$\text{and } K(m, l) \approx \frac{F_{lm}}{\Delta E_{lm} T_e^{\frac{1}{2}}} g_l/g_m \quad (2.35)$$

while the radiative decay rate (sec^{-1}) can be written in a similar form

$$A(m, l) \approx F_{lm} \Delta E_{lm}^2 g_l/g_m \quad (2.36)$$

where l , m and n are the principle quantum numbers with $l < m < n$, ΔE the energy of the transition, F_{lmn} the absorption oscillator strength, g the degeneracy of the level, K Boltzmann's constant and T_e the electron temperature. From (2.34), (2.35) and (2.36) it can be seen that collisional rates are proportional to ΔE^{-1} and radiative rates proportional to ΔE^2 , hence the dominance of collisional processes across small gaps and radiative processes across large gaps.

A population inversion can be formed in the de-excitation cascade at a point where an upper level is strongly populated from the levels above, while a lower level is rapidly depopulated to the ground state. These conditions can be met in two situations.

If there is a single large energy gap with closely spaced levels above and below, then collisional processes will dominate the transitions between closely spaced levels. However there will be a fast radiative decay across the gap, which will depopulate the levels immediately above it. Population inversions between these levels and those higher up in the collisional cascade can be formed in this way. This method has been used successfully in the optical [59] but has not shown promise in the XUV or X-ray regions.

A second method is to choose a set of levels whose spacing decreases with increasing energy. Such levels are found in systems with a single electron in the outer shell, for instance hydrogen, lithium and sodium like ions. In these systems the radiative decay rate through the cascade is larger for the lower levels, as lower energy levels have larger gaps between them. This can lead to population inversions between the lower levels depopulated by radiative decays, and the upper levels populated by collisional de-excitations from above. This method has been particularly successful in producing inversions and gain in the XUV region. Figures 2.1, 2.2 and 2.3 show simplified level diagrams for hydrogenic, lithium like and sodium like ions and the likely population inversions.

The point in the cascade at which collisional and radiative rates are approximately equal is known as the thermal limit [57]. For H like ions it can be written as :-

$$n_Q \approx 1.2 \times 10^2 Z^{14/17} n_e^{-2/17} \left[\frac{K T_e}{Z^2 E_H} \right]^{1/17} \exp \left[\frac{4Z^2 E_H}{17 n_Q^3 K T_e} \right] \quad (2.37)$$

where n_Q is the principle quantum number of the level nearest the thermal limit, n_e the electron density, E_H the ionisation energy of the hydrogen atom and Z the atomic number. For $n > n_Q$ transitions are collisionally dominated and the level populations approach the thermal equilibrium values. The population densities are then described by the Saha-Boltzmann equation [57] one form of which is given by equation (2.38), and as a result of this no inversion is possible in levels above n_Q .

$$\frac{N_{i+1}}{N_i} = \frac{1}{n_e} \frac{g_{i+1}}{g_i} 2 \left[\frac{m K T_e}{2\pi \hbar} \right]^{3/2} \exp \left[-\frac{\chi_i}{K T_e} \right] \quad (2.38)$$

where N_i and N_{i+1} are the relative number densities of the i th and $(i+1)$ th ionisation states, g_i and g_{i+1} the degeneracy of the states ($2n^2$ for H like ions) and χ_i the energy needed to ionise the i th ionisation state.

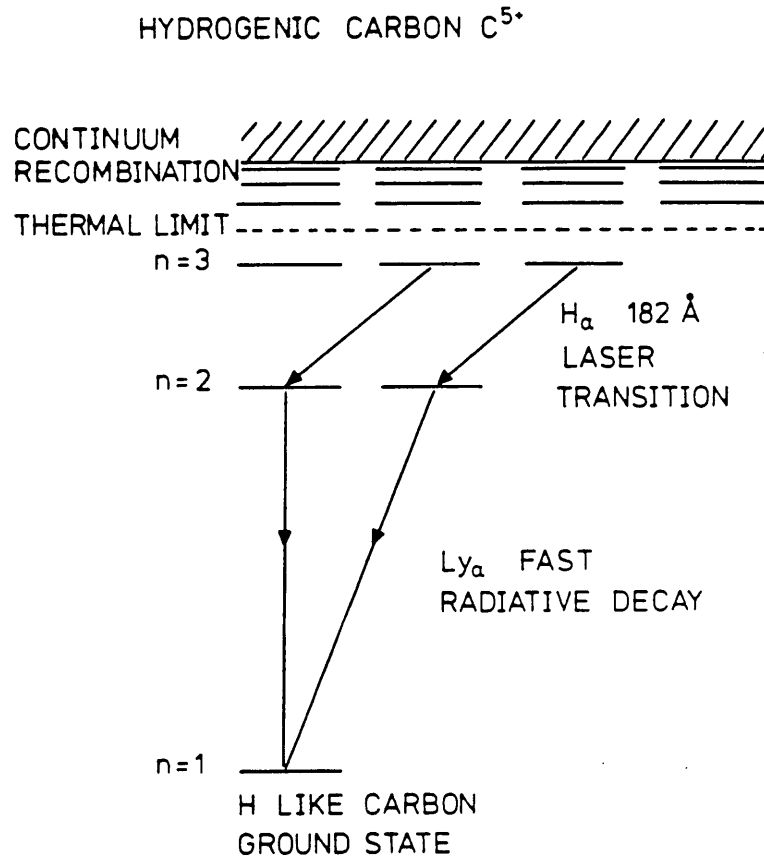


Figure 2.1 Simplified energy level diagram for the hydrogenic carbon X-ray laser scheme.

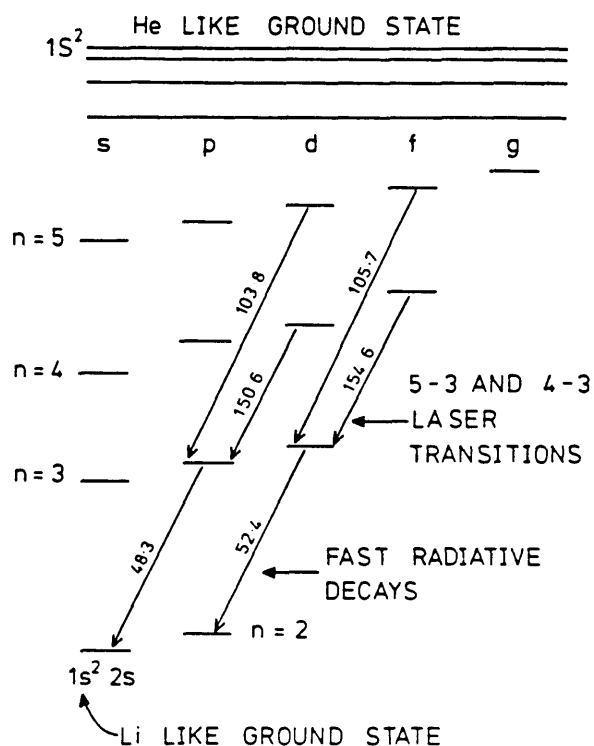


Figure 2.2 Simplified energy level diagram for the lithium like aluminium X-ray laser scheme.

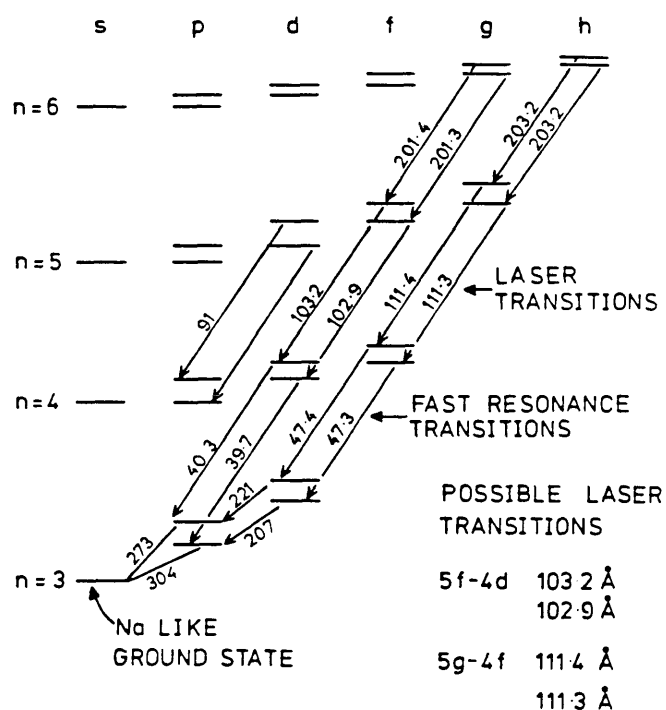


Figure 2.3 Simplified energy level diagram for the sodium like copper X-ray laser scheme.

However below n_Q where radiative processes dominate, population densities can be higher than expected simply from the radiative transition rates, due to cascades down from the collisionally dominated levels. This is the regime in which population inversions can occur. A four level laser system can operate, with the highest level 4 above the thermal limit acting as an electron reservoir to populate the upper laser level 3 below the thermal limit. The 2 level is depopulated to the ground state by a fast radiative 2–1 resonance transition (Ly_α for H like ions) leading to a 3–2 inversion.

2.6 PLASMA CONDITIONS FOR GAIN

A number of conditions must be met if significant XUV gain is to be attained in a recombining plasma. Before the plasma begins to recombine it must be ionised beyond the lasing species, fully stripped for H like schemes, He like for Li like schemes and so on. The recombination mechanism must preferentially populate levels above the laser transition rather than allowing recombination direct to the ground state. From equation (2.29) it follows that high population inversion densities are required for high gain, and consequently the upper laser level should be strongly populated. Similarly the lower laser level should be rapidly depopulated to the ground state, and there should be little repopulation of this level by reabsorption of the lower level decay radiation. The thermal limit as given by (2.37) must be above the upper laser level or collisional mixing will destroy the inversion. Finally, for a scheme with no mirrors, equation (2.33) shows that a high aspect ratio plasma is needed to give gain along one preferred axis, as the final output is a strong exponential function of the path length through the gain medium.

The work of Pert and others [46,47,31,38] has extended the recombination laser scheme into the XUV region. It has been shown that these conditions can be met in plasmas produced by laser irradiation of thin fibre targets. A high aspect ratio plasma is established by the initial target geometry, typically 10 mm long and $< 10 \mu\text{m}$ in diameter. Short pulse ($< 200 \text{ ps}$) high intensity irradiation of the fibre in line focus geometry produces a hot (few hundred eV) fully ionised plasma on a timescale short compared to that on which significant expansion takes place.

The plasma then expands rapidly over a few nanoseconds and the electron density falls from $\approx 3 \times 10^{21} \text{ cm}^{-3}$ to $\approx 2.5 \times 10^{19} \text{ cm}^{-3}$ resulting in strong adiabatic cooling. This fast cooling "freezes" the ionisation stage and allows three body recombination to preferentially populate the higher lying energy levels.

From (2.37) it can be seen that the thermal limit $n_Q \propto (T_e/n_e^2)^{1/17}$ and the thermal limit rises as the plasma expands and cools. If the thermal limit rises to just above the upper laser level a population inversion can form. The upper lasing level will be in approximately Saha-Boltzmann equilibrium with the higher lying levels and strongly populated by collisional cascades down from these levels. For hydrogenic ions the $n = 3$ upper laser level population N_3 is given by :-

$$N_3 = \frac{n_e n_i}{T_e^{3/2}} \exp \left[\frac{Z^2 E_H}{9 K T_e} \right] \quad (2.39)$$

Where n_e is the electron density and $n_i \approx n_e/Z$ the bare ion density. This shows that high electron densities and low temperatures are required for strong population of the upper laser level. However the lower laser level must also lie well below the thermal limit or collisional mixing will destroy the population inversion. If the lower level is below the thermal limit then it will be depopulated by a rapid radiative transition to the ground state. To avoid repopulation of the lower laser level by reabsorption of this radiation, the plasma must be optically thin for the lower level to ground state transition. The optical depth τ for radiation escape follows from equations (2.29) and (2.32) and can be written as :-

$$\tau = \lambda^2 g(\nu) A_{21} N_1 R \quad (2.40)$$

where N_1 is the ground state population density and R the path length through the plasma.

The expanding cylindrical plasma produced by laser irradiation of thin fibre targets minimises radiation trapping in three ways. Typically the gain in such a scheme will turn on and decay in under two nanoseconds, due to the precision with which n_e , n_i and T_e must be matched. In this short time the hydrodynamic expansion of the plasma in the radial direction is only on the order of a few hundred microns and thus R the path length is small.

Secondly the plasma is initially fully stripped beyond the lasing species, and since fast cooling avoids direct recombination to the ground state, the ground state population density N_1 is low.

Finally the small initial diameter of the fibre and high initial density and temperature of the plasma results in a rapid radial expansion. This expansion gives rise to the strong adiabatic cooling needed to establish the recombination cascade. However it also results in a radial velocity profile with a steep gradient. The radial expansion of the plasma Doppler shifts the lower laser level to ground state decay radiation, but due to the velocity gradient, the Doppler shift is different for different radial components of the plasma. This differential Doppler shifting in the radial direction reduces the line overlap for escaping radiation. As a result of these three effects, the absorption length can be much smaller than the radial dimension of the plasma, and the plasma can become optically thin in the radial direction.

2.7 OPTIMISATION OF PLASMA CONDITIONS

Clearly there will be a set or sets of plasma conditions which will maximise gain in a fibre based recombination laser. These conditions have been examined in detail, particularly by Pert who used self-similar expansion models [47] coupled to collisional-radiative rate equations to identify expansion adiabats for which the time evolution of population, temperature and electron density lead to peak gain. This also led to a set of semi-empirical scaling laws for the gain in terms of experimental parameters.

The most important of these parameters are the target dimensions and atomic number, and the pump laser pulse length and irradiance which control the mass, energy density and expansion profile of the ablated plasma. These scaling laws can be used both to design experiments and to scale working schemes to shorter wavelengths by changing the target material and experimental conditions. Perts original work on hydrogenic carbon plasmas has since been generalised to a range of higher Z elements. Starting with a hot, uniform cylinder of fully ionised (carbon) plasma of mass M (g cm^{-1}) per unit length and radius R (cm) the optimum thermal energy per unit length identified by simulations is :-

$$E_t = 1.9 \times 10^{11} M^{7/4} R^{-3/4} \quad (\text{J}) \quad (2.41)$$

The corresponding peak gain coefficient α is :-

$$\alpha = 3.9 \times 10^{-6} M^{-1/2} R^{-1} \quad (\text{cm}^{-1}) \quad (2.42)$$

The $M^{-1/2}$ scaling of α is at first glance unphysical, more plasma gives less gain. However this follows from the balance between inversion density which from (2.39) is approximately proportional to n_e^2 and the opacity for the lower state decay radiation Ly_α . Clearly (2.42) will fail in the very low mass limit.

These results do not take into account the mechanism by which the plasma cylinder is produced. The time evolution of the plasma can be described by two essentially separate phases. In the first "heating phase" laser ablation of the target establishes the dimensions, mass and energy density of the plasma on a timescale short compared to the expansion time of the plasma. These must match the conditions needed to put the plasma on an adiabat which leads to gain.

The second phase describes the subsequent expansion and cooling of the plasma, and it is this "expansion phase" which must satisfy (2.41) and (2.42). The radius of the plasma cylinder in (2.41) can be identified with the radius of the plasma established by hydrodynamic expansion during the initial heating phase.

$$R = t (E_t/M)^{1/2} \quad (2.43)$$

Where t is determined by the laser pulse shape and duration, typically the time to peak. The optimum expansion adiabat for gain is then described by :-

$$E_t = 1.3 \times 10^6 M^{17/11} t^{-6/11} \quad (2.44)$$

$$\alpha = 1.6 \times 10^{-12} M^{-17/22} t^{-8/11} \quad (2.45)$$

As with (2.42) the equation for gain (2.45) fails in the very low mass limit where the expansion time becomes less than the heating time. The total energy absorbed by the plasma during the heating phase is the sum of the ionisation and thermal energies $E_{\text{tot}} = E_i + E_t$:-

$$E_{\text{tot}} = 8.3 \times 10^6 M + 1.3 \times 10^6 M^{17/11} t^{-6/11} \quad (2.46)$$

2.8 LASER ABLATION

The mass of the plasma cylinder is determined by the ablation mechanism and the total energy E_{tot} absorbed during the heating phase. The conditions for maximum gain also require that the absorbed energy be a function of the plasma mass ablated M_b , thus :-

$$M_b = M_b(E_{\text{tot}}) \quad (2.47)$$

$$E_{\text{tot}} = E_{\text{tot}}(M_b) \quad (2.48)$$

and hence the optimum parameters are described by

$$E_{\text{tot}} = E_{\text{tot}}(M_b(E_{\text{tot}})) \quad (2.49)$$

The relationship between M_b , E_{tot} and α is not immediately obvious from (2.43), (2.44) and (2.45), and it is most easily seen on a "burn diagram". Figure (2.4) is an idealised burn diagram due to Pert [47] which plots M/M^* against E/E^* where M^* and E^* are the optimum values of plasma mass and energy given by equations (2.45) and (2.46) for peak gain, and the point of intersection with the gain curve indicates the optimum conditions. From the plot it can be seen that there are two operating regimes, one at high absorbed energy and high mass and one at low mass and low absorbed energy. From equation (2.45) it follows that since $\alpha \propto M^{-1/2}$ the low mass regime will give higher gain.

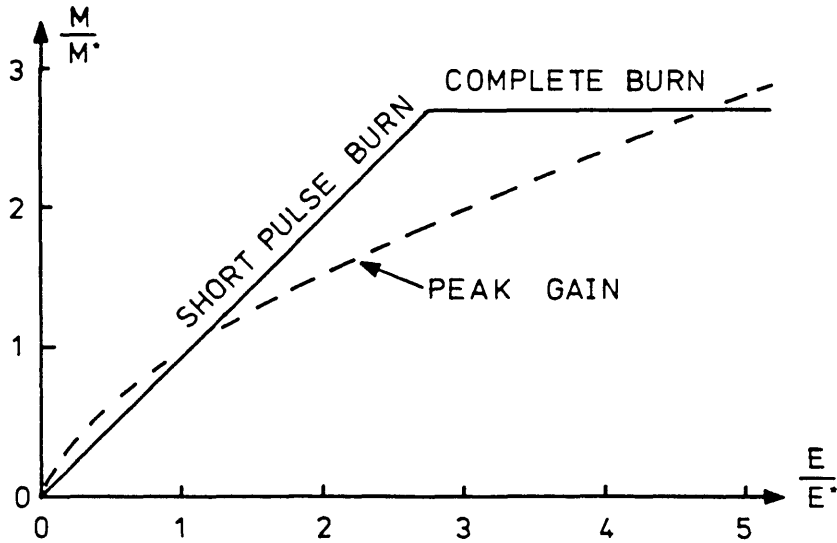


Figure 2.4 Idealised burn diagram showing the relationship between optimum plasma mass M^* and optimum plasma energy E^* .

Experimentally the plasma is produced by irradiating a thin fibre target with 70 – 200 ps pulses of green ($\lambda = 0.53 \mu\text{m}$) laser light at intensities in the region of $10^{14} \text{ W cm}^{-2}$. In this regime there is sufficient time for a significant underdense plasma to form during the heating pulse, and as a result the laser absorption mechanism is predominantly self-regulating inverse Bremsstrahlung [47]. The mass of plasma ablated is :—

$$M_b = 2.6 \times 10^{-3} \gamma Z_i^{-1} A_m^{7/9} \lambda^{-4/9} r_0^{2/9} E^{5/9} t^{4/9} \quad (2.50)$$

where γ is a laser pulse shape factor ($\gamma \approx 1.18$ for a Gaussian), λ (μm) and t (sec) the pump laser wavelength and pulse width, Z_i the ionisation number of the plasma, A_m and r_0 (cm) the target mass number and radius, and E (J cm^{-1}) the absorbed energy per unit length. Substituting (2.50) into the expression for gain (2.45) then gives :—

$$\alpha \approx 7.3 \times 10^{-9} \gamma^{-0.77} Z_i^{0.77} A_m^{-0.6} \lambda^{0.34} r_0^{-0.17} E^{-4.3} t^{-1.1} \quad (2.51)$$

Note that for a given target material smaller diameter targets and shorter pump laser pulse lengths give higher gain. However the target diameters cannot be made arbitrarily small as they are limited to $\geq 7 \mu\text{m}$ by experimental stability requirements. In the high gain, low mass regime only a thin layer of the target ($\leq 1 \mu\text{m}$) is generally ablated, and this results in a hot plasma surrounding a cold solid core. Laser energy is absorbed in the underdense plasma at some distance from the core, and this results in a nearly isothermal annular plasma which rapidly expands away from the core after the heating phase, and becomes decoupled from it [60].

2.9 SHORT WAVELENGTH SCALING

One of the advantages of the hydrogenic recombination XUV laser scheme is its ease of scalability to shorter wavelengths by moving to higher Z lasing ions. The simple Bohr model of the hydrogen like ion gives a strong Z^{-2} scaling of the wavelength of the Balmer H_α 3–2 laser transition with atomic number. The 3–2 laser transition wavelengths are tabulated for a range of low Z elements. The wavelengths for these transitions recorded in the table on page 53 do not include the fine structure splitting of the lines. The fine structure components cannot be resolved with our present diagnostics, however their effect on the predicted gain is included in detailed simulations, and is discussed later.

Element	Z	Wavelength Å
C	6	182
N	7	134
O	8	102
F	9	81
Ne	10	66
Na	11	54
Mg	12	46
Al	13	39

However the conditions outlined in sections 2.5 – 2.8 must still hold for there to be significant gain. The population inversion in hydrogenic recombination schemes arises from the balance between collisional and radiative de-excitation rates above and below the thermal limit. Clearly if the scheme is scaled to higher atomic numbers the thermal limit must still remain above the $n = 3$ upper laser level for there to be any chance of a 3–2 inversion. The position of the thermal limit n_Q is given by (2.37) and is determined by the ratio of collisional to radiative transitions. For n_Q to remain approximately in the same place, the ratio of collisional to radiative transitions must be constant for a range of atomic numbers Z .

Examining the expressions for the rate coefficients (2.34 – 2.36) we find that over a limited range where the $\exp(-\Delta E_{mn})$ dependence of the $m \rightarrow n$ collisional rates can be ignored :–

$$\text{Collisional rates} \propto n_e / \Delta E \cdot T_e^{\frac{1}{2}} \quad (2.52)$$

$$\text{Radiative rates} \propto \Delta E^2 \quad (2.53)$$

For hydrogenic ions the energy levels scale as Z^2 and so $\Delta E \propto Z^2$. To maintain the ratio of collisional to radiative transitions :–

$$\frac{\text{Collisional Rate}}{\text{Radiative Rate}} = \frac{n_e}{\Delta E^3 T_e^{\frac{1}{2}}} = \text{constant} \quad (2.54)$$

Clearly from the expression for the thermal limit (2.37) T_e must scale as Z^2 , and since $\Delta E \propto Z^2$ the density must go as:–

$$n_e / Z^7 = \text{constant} \quad (2.55)$$

While the collisional and radiative rates used to derive the $T_e \propto Z^2$ and $n_e \propto Z^7$ relationships are not exact, they are valid over a wide range of conditions [61]. The Z scaling of the gain can now be found from equation (2.29) which gives $\alpha(\nu) = \lambda^2/8\pi g(\nu) A_{21} \Delta N$. Since the laser ion is initially fully ionised and N_2 is small we can set $\Delta N \approx N_3$ the upper laser level population density. N_3 is given by a modified form of the Saha-Boltzmann equation (2.39), while $g(\nu)$ the line profile is predominantly due to Doppler broadening and is given by (2.20). Substituting our expressions for N_3 and $g(\nu)$ into (2.29) gives the gain as :-

$$\alpha(\nu) = \frac{\lambda^2}{8\pi} \frac{\lambda}{\pi} \left[\frac{m \pi}{2 K T_e} \right]^{\frac{1}{2}} A_{21} \frac{n_e n_i}{T_e^{3/2}} \exp \left[\frac{Z^2 E_H}{9 K T_e} \right] \quad (2.56)$$

We already have $\lambda \propto Z^{-2}$, $T_e \propto Z^2$ and $n_e \propto Z^7$. The ion density $n_i \approx n_e/Z$ thus $n_i \propto Z^6$. $K T_e$ will typically be large compared to $Z^2 E_H$ since the plasma is fully stripped and so we may ignore the exponential term in the expression for N_3 and find that $N_3 \propto Z^{10}$. The ion mass scales approximately as $m \propto Z$ and hence $g(\nu) \propto Z^{-5/2}$. For hydrogenic ions the absorption oscillator strength $A_{21} \propto Z^2/\lambda \propto Z^4$ [57] and so $\alpha(\nu) \propto Z^{-4} \cdot Z^{-5/2} \cdot Z^4 \cdot Z^{10} = Z^{7.5}$ which looks very favourable. Extensive computer modelling by Pert for a range of low to medium Z hydrogenic schemes gives the optimum scaling of parameters for maximum gain as :-

Wavelength	λ	$\propto$	Z^{-2}
Energy	E	$\propto$	Z^2
Gain	α	$\propto$	$Z^{7.5}$
Temperature	T_e	$\propto$	Z^2
Density	n_e	$\propto$	Z^7
Time	t	$\propto$	Z^{-4}

(2.57)

However it may not be possible to achieve all these conditions experimentally. The electron temperature T_e of a plasma produced by inverse Bremsstrahlung absorption of laser light is related to the irradiance I by $T_e \propto Z^{2/9} I^{4/9}$ [62]. To give the required $T_e \propto Z^2$ scaling for optimum gain the pump laser intensity must therefore scale as $I \propto Z^4$, which is limited by the available pump power rather than any physical processes.

The Z^7 scaling of initial density needed to achieve optimum gain is more of a problem. Laser ablation produces a plasma whose initial density scales as $n_e \propto Z$ and this results in a deviation from the ideal parameters (2.57). Gain is only possible when the thermal limit is above the $n = 3$ level, and this condition is density dependent. The initial $n_e \propto Z$ value for the laser produced plasma must converge to the ideal $n_e \propto Z^7$ value as the plasma expands and cools. This is achieved by decreasing the plasma expansion ratio and the associated adiabatic cooling with increasing atomic number. The expansion ratio for the density in terms of the initial n_i and final n_f densities is :-

$$\frac{n_f}{n_i} \approx \left[\frac{R_i}{R_f} \right]^2 \quad (2.58)$$

and the corresponding cooling ratio is :-

$$\frac{T_f}{T_i} \approx \left[\frac{R_i}{R_f} \right]^{4/3} \quad (2.59)$$

To give the optimum final density scaling $n_f \propto Z^7$ requires $n_i (R_i/R_f)^2 \propto Z^7$ and hence $(R_i/R_f) \propto Z^3$. Substituting $T_i \propto Z^2$ and $(R_i/R_f) \propto Z^3$ into (2.58) gives a $T_f \propto Z^6$ scaling of the final electron temperature after the expansion. These experimentally achievable scalings of n_e and T_e may be substituted back into (2.56) to find the modified scaling of the gain with atomic number. The population inversion density now scales as $\Delta N \approx N_3 \propto Z^4$ and $g(\nu)$ as $Z^{-9/5}$ to give a gain scaling of $\alpha \propto Z^{-1/2}$ which is not as favourable as the optimum $Z^{7.5}$, but still results in significant gain for higher Z lasing ions at shorter wavelengths. The modified initial conditions then scale as :-

Wavelength	λ	$\propto$	Z^{-2}
Energy	E	$\propto$	Z^2
Gain	α	$\propto$	$Z^{-1/2}$
Temperature	T_e	$\propto$	Z^2
Density	n_e	$\propto$	Z
Irradiance	I	$\propto$	Z^4
Time	t	$\propto$	Z^{-1}

(2.60)

Computer simulations show that significant gain is possible with the hydrogenic recombination scheme up to Z values of 12 – 13 and at the associated wavelengths of 46 – 39 Å. Chapter 5 describes the successful experimental scaling of this scheme from carbon ($Z = 6$, $\lambda = 182$ Å) to fluorine ($Z = 9$, $\lambda = 81$ Å). The limit on the short wavelength scalability of this scheme is imposed by the initial density which approaches solid for aluminium ($Z = 13$).

Solid density plasmas are difficult to produce in geometries other than spherical compression if the pump laser pulse length is > 20 ps. An experiment describing the first production of fully stripped, near solid density Al plasmas in slab geometry using picosecond pulses is reported in chapter eight.

The high initial density requirement for gain may be relaxed if the plasma can be cooled by other means in addition to adiabatic expansion. Radiation loss in particular may enhance the cooling, as radiation loss due to Bremsstrahlung and radiative recombination scales as Z^2 while line radiation loss scales as Z^4 . Doping XUV laser targets with small amounts of high Z elements may result in significant additional cooling and this is discussed in chapter seven.

CHAPTER 3. EXPERIMENTAL FACILITIES AND DIAGNOSTICS

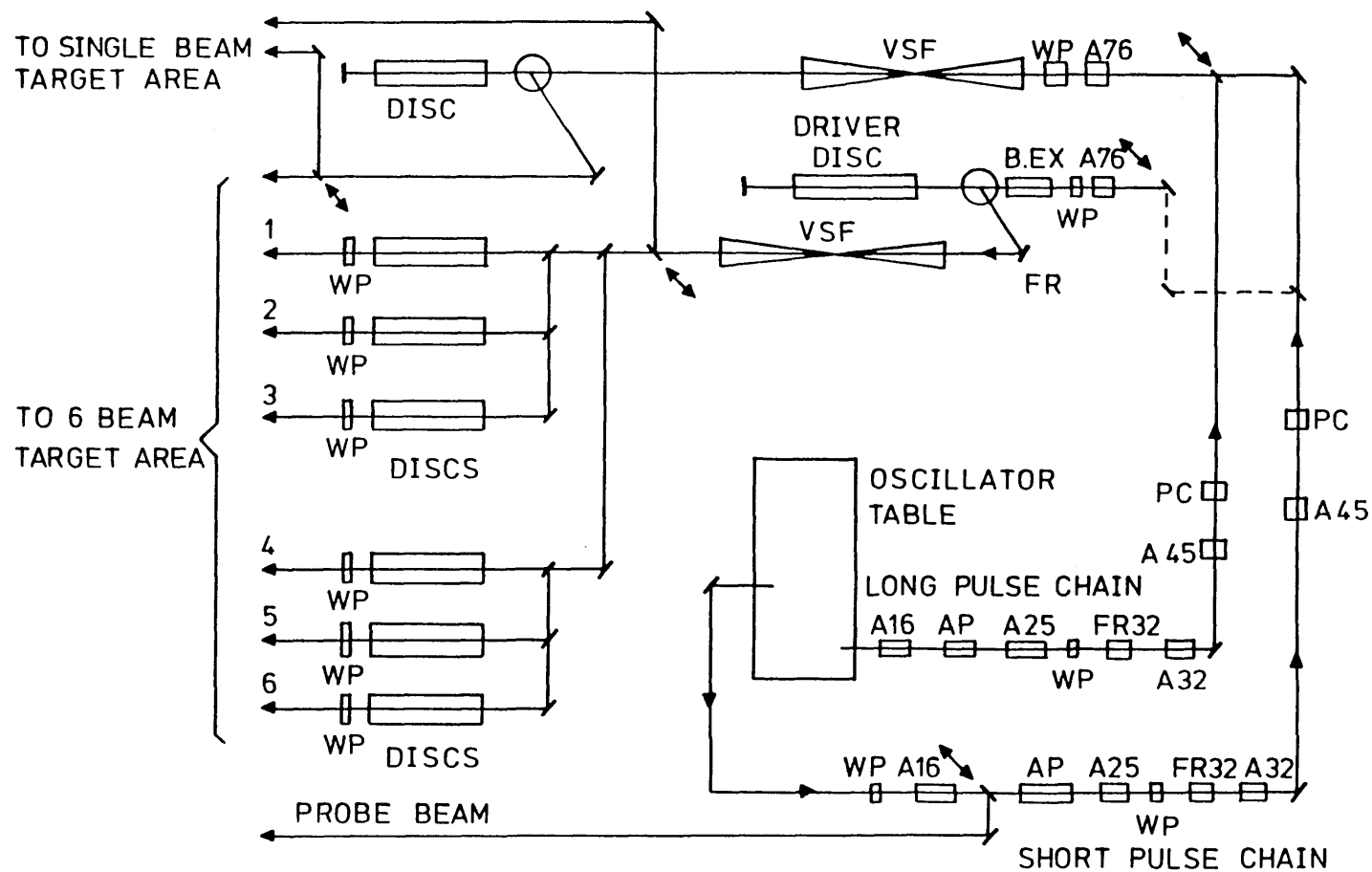
The majority of the experiments described in this thesis were carried out at the Science and Engineering Research Council's (SERC) Rutherford Appleton Laboratory (RAL). This Laboratory provides large facilities and technical support for visiting university research groups. The various X-ray laser schemes studied were powered by the VULCAN neodymium glass laser system [63]. A novel six beam off axis lens / mirror line focus [64] was used to focus up to 2 TW (150 J, 70 ps) pulses of frequency doubled green ($\lambda = 0.53 \mu\text{m}$) laser light onto thin fiber and foil X-ray laser targets. The plasmas produced were studied with a number of state of the art diagnostics. These included time resolved flat field XUV spectrometers developed specifically for X-ray laser research, and a range of X-ray pinhole cameras, plasma calorimeters, Faraday cups and X-ray crystal spectrometers originally used in laser fusion and laser plasma interaction experiments.

3.1 THE VULCAN GLASS LASER SYSTEM

VULCAN is an infra-red ($\lambda = 1.053 \mu\text{m}$) multi-beam high power ($\approx 3 \text{ TW}$) Nd:Glass laser system, and was originally developed for laser fusion studies. However this type of laser can also deliver the large pump powers on short timescales required by X-ray laser experiments.

The VULCAN system consists of synchronous long (0.7 – 20 ns FWHM) and short (70 – 500 ps FWHM) oscillators which provide 200 μJ seed pulses for a series of Nd:Phosphate glass amplifiers (see figure 3.1). For the reasons outlined in chapter 2, most experiments were carried out with short pulses in the 70 – 200 ps range. The short pulse oscillator is a phosphate glass system operating at $\lambda = 1.053 \mu\text{m}$, actively mode locked by an acousto-optic loss modulator and Q switched with an electro-optic Pockels cell. The pulse duration can be varied by using different intra-cavity etalons. A single pulse from the oscillator is switched out and passed through a diffraction limited spatial filter which injects a 5 mJ expanding quasi-Gaussian beam into the amplifier chain. The spatial distribution of the beam is modified by an apodizer and the radial gain profile of the rod amplifiers to provide as large a fill factor as possible without producing severe edge diffraction effects.

Figure 3.1 The VULCAN multi Tera-Watt Nd:Glass laser system.



The pulse passes through a series of Xe flashlamp pumped rod amplifiers of expanding diameter to avoid non linear effects and damage to optics by keeping the peak power / unit area below 3 J cm^{-2} . A series of wave plates and Faraday rotators protect the system from back reflections off components further up the amplifier chain.

At the end of the rod amplifier chain the $1 - 1.5 \text{ J}$ pulse is spatially filtered to improve beam quality and double passed through a 108 mm diameter disc amplifier. The beam is then split six ways and the six beams single passed through large aperture 108 mm disc amplifier modules. The final output of the laser is $\approx 300 \text{ J}$ total infra-red in six 50 J beams. This is a maximum output for 70 ps pulses imposed by non linear effects in the final amplifiers and by the damage threshold of the dielectric multi-layer mirrors used to steer the beams. The output can be varied downwards in energy by changing the gain staging of the amplifier chain, and the beam to beam and shot to shot variation in output energy is $\approx 20 \%$.

The six beams pass into a switchyard which can route the output of the laser to several target areas. The beams are frequency doubled to $\lambda = 0.53 \text{ }\mu\text{m}$ in large aperture 15 mm thick type II KDP crystals which give $\approx 40 - 50 \%$ conversion efficiency when tuned correctly. Green ($\lambda = 0.53 \text{ }\mu\text{m}$) light rather than the $1.053 \text{ }\mu\text{m}$ first harmonic output of the laser is used as this gives a greater coupling efficiency to X-ray laser targets. Each beam is passed along a variable length "trombone" which can delay individual beams by several nanoseconds. The trombones are used to ensure that the six beams which all travel slightly different routes through the laser and target area arrive on target coincident in time. Calorimeters monitor the infra-red and green energy of each beam, both for crystal tuning and to give the total energy on target.

3.2 LINE FOCUS OPTICS

One major problem for previous workers in the field has been how to efficiently couple pump laser energy into a long, thin X-ray laser target. A novel six beam system which uses off axis lens / mirror pairs [64] rather than cylindrical lenses to produce a high quality diffraction limited line focus [65] was used to perform the X-ray laser experiments described in this thesis.

Any optical system which has rotational symmetry about one axis, illuminated by a diffraction limited source on the axis of symmetry will give a diffraction limited line focus if there is a real image. The line focus will also lie along this axis of symmetry. This assertion is based purely on symmetry and the fact that all refracted or reflected rays must lie in the plane of incidence. A number of optical systems will satisfy the rotational symmetry criteria, and the one used in our case was an F / 2.5 aspheric doublet lens [66] and spherical mirror (see figure 3.2). Both of these components can be made to high tolerances and are capable of producing near diffraction limited foci.

The length of the line focus produced by the lens / mirror pair is a function of the mirror tilt angle θ and the F number. For a given θ and F all distances scale linearly with the mirror radius R . The line focus length L is given by :-

$$L = 1.4 \times 10^{-4} R \frac{\theta^2}{F \cos \theta} \quad (3.1)$$

or more simply :-

$$L = 1.4 \times 10^{-4} \theta^2 D \quad (3.2)$$

where D is the illuminated length on the mirror and θ is the tilt angle in degrees. From equation (3.2) it is clear that long line foci in the order of several cm can be produced by this method. For an unaberrated input beam the width w of the line focus is given by :-

$$w = 2.4 F \lambda \cos \theta = 2.4 \lambda (R/D) \quad (3.3)$$

while for an input beam with divergence β and a primary lens focal length f :-

$$w = \beta F \cos \theta = \beta (R/D) \quad (3.4)$$

Note that in both (3.3) and (3.4) the width of the line focus decreases with increasing θ , while the line focus length increases. Even for an input beam with greater than 100μ rad divergence it is still possible to achieve line focus aspect ratios of $> 10^4$ as compared to $\approx 10^2$ for a conventional cylindrical lens.

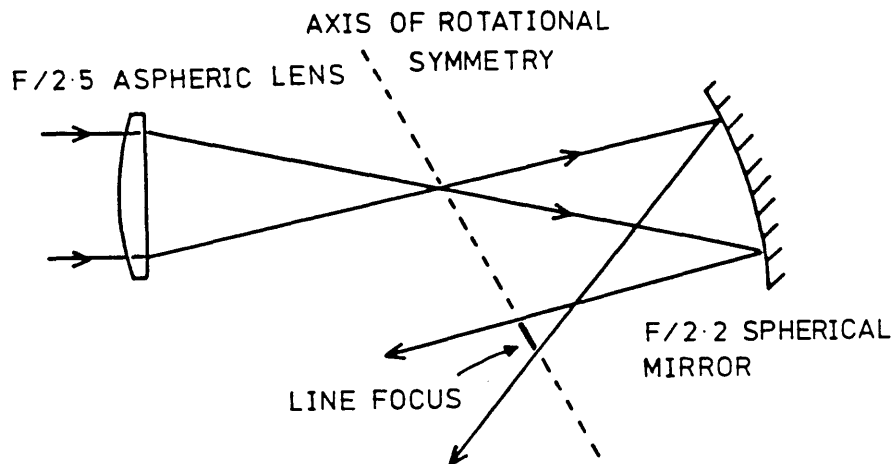


Figure 3.2 The off-axis mirror technique for producing a line focus.

The TA East line focus facility at the Rutherford Appleton Laboratory uses six off axis lens / mirror pairs to build up one long line focus. Each pair comprises an $F / 2.5$ input lens and 280 mm diameter spherical mirror with a tilt angle of 26° . An input beam of ≈ 80 mm in diameter gives a 7.5 mm long line focus, limited by the 100μ rad divergence of the VULCAN laser to 25μ m in width. The six beams are arranged in three opposing pairs at 30° to each other (see figure 3.3). These pairs may be superimposed to give high intensity, symmetrical irradiation over lengths up to 7 mm, or the pairs may be spaced out along a target to give lower intensity irradiation over lengths up to 21 mm. The performance of a single segment of this line focus has been examined experimentally by previous workers [67]. Plasma calorimeter arrays indicated coupling efficiencies of 10 – 15 % with 10μ m fibre targets, and pinhole camera images showed the resulting plasmas to be very uniform over their full length.

The full six beam facility has since been used on a number of highly successful X-ray laser experiments [chapters 4,5,6] in which XUV gain was demonstrated in several different recombination schemes.

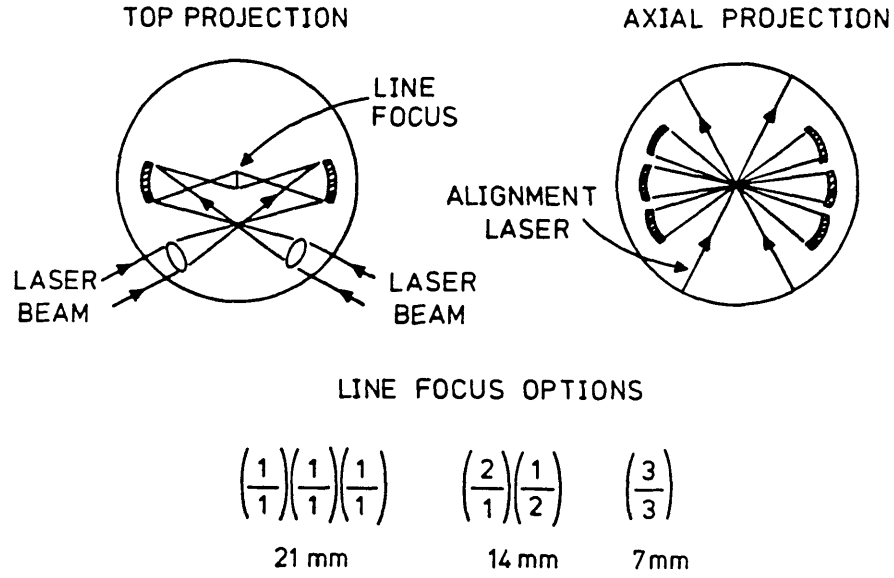


Figure 3.3 Variable length line focus production by three opposing beam pairs.

3.3 LINE FOCUS AND TARGET ALIGNMENT

At the start of each experiment the line focus was set up using a 20 μm diameter tungsten wire surrogate target strung along the chamber axis to define both the line focus, and the target position in the alignment optics. The six segments of the line focus were then aligned onto this wire with a CW argon ion laser ($\lambda = 0.528 \mu\text{m}$) which has a similar wavelength to the 0.5265 μm second harmonic of the VULCAN laser, and which was colinear with the VULCAN beams. An optical streak camera was used to image a 1 mm steel ball placed at the primary focal point of the six aspheric doublet lenses. Low power VULCAN shots were fired, and the arrival times of the six beams adjusted by varying their path lengths such that they all arrived coincident in time to better than 15 picoseconds.

The divergence of a mirrorless ASE laser ≈ 10 mm long and 200 μm wide is on the order of 20 milliradians. A typical target is initially some 10 μm in diameter and must be uniformly irradiated by a 25 μm wide line focus. This gives a rough feel for the accuracy to which targets must be aligned, both to be hit by the pump laser, and for the diagnostics to be in the correct position to detect any ASE signal from the plasma.

Targets were aligned under vacuum using an imaging system comprising a pair of orthogonal split field microscopes [68]. This system was set up on the same wire target used to establish the line focus position. Using a motorised target mount with three linear and two angular degrees of freedom, targets could be repeatably positioned with an accuracy of better than 1 milliradian in angle and a few microns in position.

For a target shot the six segments of the line focus were then individually aligned on target. This was done by observing the "Foucault" shadow in a CW argon ion alignment beam transmitted past the target, and maximising the shadow. Motorised drives on each lens and mirror allowed this to be done under vacuum, and so no corrections for refraction of the main beams in air were needed. One unique feature of this line focus is that the alignment procedure does not effect the line focus quality, as it only tilts or translates the line focus without introducing any transverse aberrations. For shots of less than the full line focus length, some of each beam was then masked off at the input to the target chamber.

3.4 TIME RESOLVED FLAT FIELD SPECTROMETER

The high energy densities required by X-ray lasers and the consequent rapid expansion and cooling of laser plasmas results in strong XUV emission over timescales of a few tens of nanoseconds. The time over which there is a measurable ASE signal may be much shorter than this due to the precision with which temperature, density and expansion velocity must be matched for there to be significant gain. Timescales on the order of a few hundred picoseconds for the turn on and decay of gain are expected for recombination XUV lasers [chapter 4, 5]. For this reason it is necessary to have time resolved diagnostics which can distinguish a short lived but intense ASE signal from both the longer lived spontaneous background plasma emission, and the strong Bremsstrahlung continuum emitted by the plasma during the laser heating phase. The time evolution of the laser line can also be a good indicator of gain, and can be useful in testing the accuracy of numerical simulations.

Grazing incidence Rowland circle spectrometers [69] have rigidly fixed source, grating and image positions which lie on an arc of a circle. A curved image plane makes this type of instrument difficult to couple to active or time resolving detectors.

However "flat field" gratings for the XUV have recently become available [12,13]. As with Rowland circle instruments, the grating has a fixed source and image position, but the gratings periodicity varies across its surface in such a way as to distort the image onto a flat plane. This type of grating has been coupled to soft X-ray streak cameras to give the high temporal resolution which XUV spectrometers for X-ray laser research require [14].

A flat field time resolved spectrometer was constructed by the author to compliment a similar instrument used at RAL, allowing the first simultaneous axial and transverse time resolved spectroscopy of X-ray laser experiments to be carried out. These instruments were also absolutely calibrated by the author and co-workers using the Synchrotron Radiation Source (SRS) at Daresbury.

The instrument comprises an entrance slit or pinhole assembly, filter slide, and a Hitachi 1200 or 2400 lines / mm mechanically ruled, blazed and gold coated flat field grating (see figures 3.4 and 3.5). The grating is held in a motorised precision mount with two translational and two rotational degrees of freedom, allowing it to be aligned under vacuum. The grating disperses a 30 – 350 Å (1200 lines / mm) or 10 – 250 Å (2400 lines / mm) spectrum over a 50 mm flat field with resolution $\Delta\lambda/\lambda \approx 10^{-3}$. For time integrated spectroscopy an XUV phosphor plate is used to record the full spectral range, with readout onto film via a fiberoptic coupler and image intensifier.

The phosphor plate can be replaced by a soft X-ray streak camera [70] which gives continuous time resolution of a 1D section of the flat field over a 100 Å window. The spectral window is limited by the 20 mm active length of the streak camera photocathode, but the camera can be easily moved to cover any 100 Å section of the full 300 Å range of the instrument. With a streak camera detector and low density CsI photocathode the spectral resolution is reduced to $\Delta\lambda/\lambda \approx 5 \times 10^{-3}$ by the mesh size of the photocathode support structure and the grain size of the image intensifier readout.

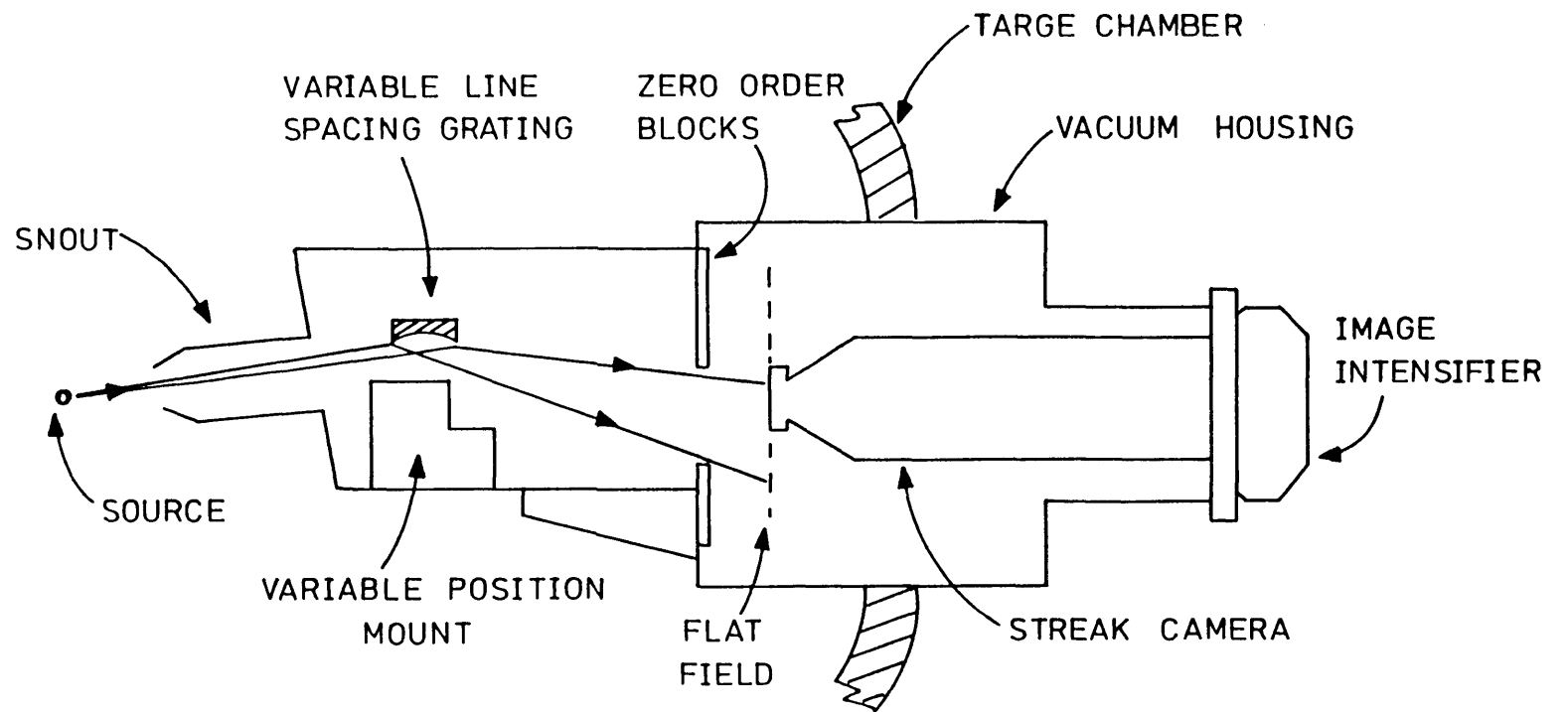


Figure 3.4 Schematic of the time resolved flat field XUV spectrometer.

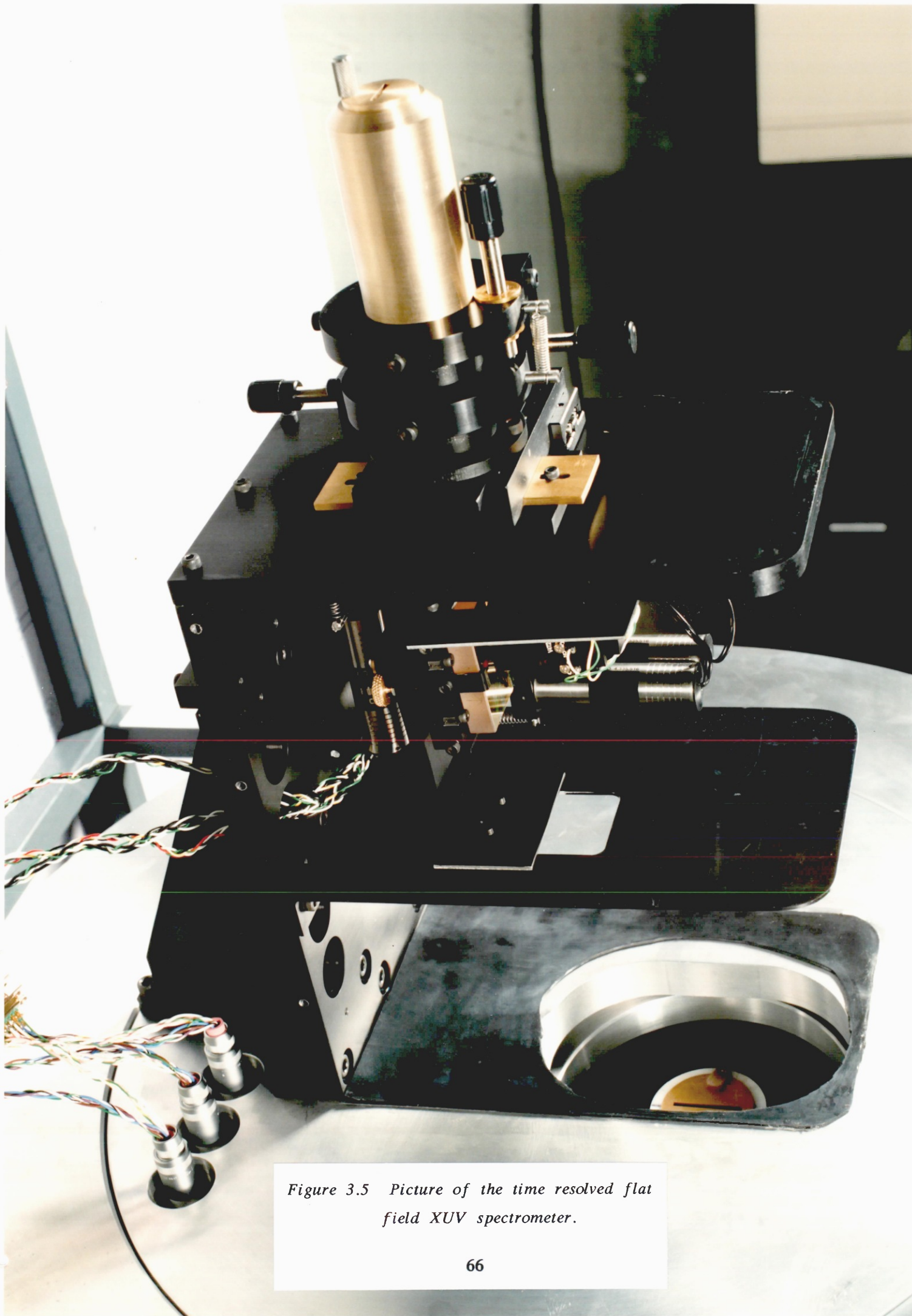


Figure 3.5 Picture of the time resolved flat field XUV spectrometer.

The temporal resolution of the spectrometer on an X-ray laser experiment is typically better than 80 ps over a 3 ns window. The finite resolution is due to the width of the streak camera input slit giving some temporal smearing of the final image. The absolute temporal resolution of this type of instrument is ≈ 15 ps when used on the fastest (15 ps / mm) sweep speed. This is caused by the energy spread of secondary photoelectrons from the streak camera photocathode. The instrument is highly sensitive in the XUV and can easily record spectra from single target shots.

3.5 SPECTROMETER ALIGNMENT

Figure 3.6 shows the flat field focusing of a Hitachi 1200 lines / mm grating for a range of input angles. This type of grating was typically used at 3° grazing incidence (87° normal incidence in figure 3.6). As can be seen in figure 3.6 the flat field rapidly defocuses if the incidence angle is varied by more than 0.5° . This becomes even more critical for the 2400 lines / mm grating which is used at 1.5° grazing incidence.

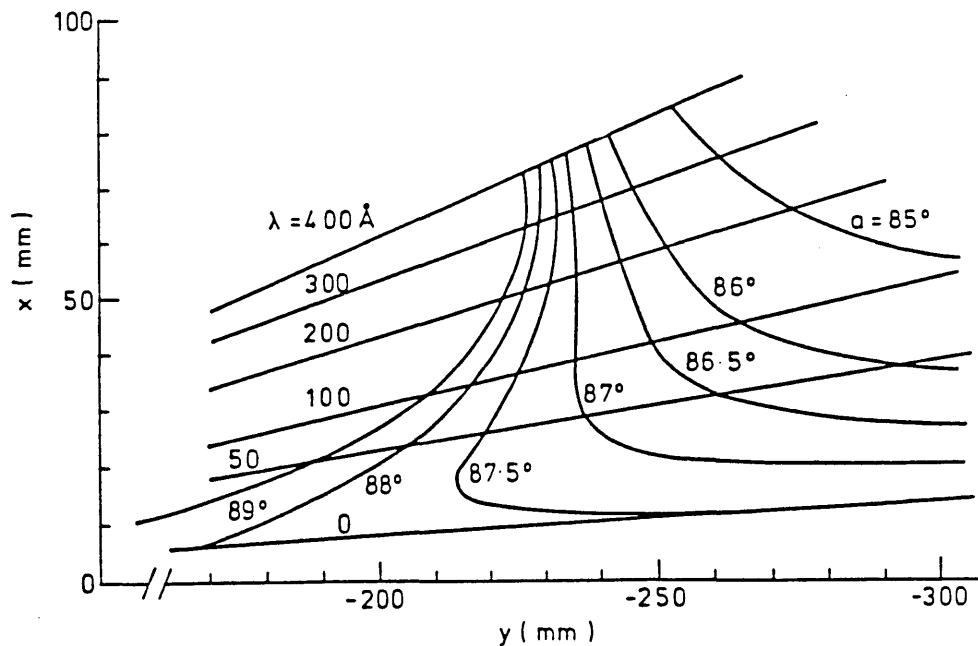


Figure 3.6 Flat field focusing of the 1200 lines / mm grating for a range of input angles.

To align the instrument, the source to grating and grating to flat field distances were bench set with reference to the spectrometer housing. With the spectrometer bolted to the target chamber, a collimated He Ne reference laser beam was then set to pass through the source and slit positions. The grating was driven out of the way of the alignment beam and the straight through path marked. The grating was then driven back down into the reference beam, and the zero order reflection off the grating used to set the grating angle with respect to the straight through beam. The grating was then adjusted in height until the reference beam was centered on the grating. Finally, zero order blocks were positioned to cut out all the straight through reflection off the grating, as the zero order is usually sufficiently bright to overload the detector.

3.6 SPECTROMETER CALIBRATION

An absolute calibration of the flat field time resolved spectrometers used on X-ray laser experiments is not required simply to demonstrate gain. However to compare the power outputs of our laser schemes to computer simulations and experiments by other workers, an accurate calibration as a function of wavelength is needed for the complete spectrometer / streak camera unit. Individual components of the spectrometer such as filters and gratings, as well as the complete unit were calibrated [91] against a National Bureau of Standards (NBS) vacuum photodiode [71] using a synchrotron light source. A flat Al cathode and cylindrical anode photodiode was used to carry out the experiments, and at the end of the run this was cross calibrated against an NBS diode consisting of an Al coated quartz cathode with enhanced ($150 \text{ \AA}$) oxide layer and stainless steel cylindrical anode. The NBS diode had been standardised at the NBS electron storage ring against a rare gas ionisation chamber.

The SRS at Daresbury is a 31 meter diameter 2 GeV electron storage ring, typically carrying a multi-bunch current of 250 mA in 156 individual "buckets". The ring produces 520 MHz pulses of light which are spectrally continuous from long wavelength radio ($\lambda = 100 \text{ m}$) to hard X-rays ($\lambda = 0.1 \text{ \AA}$). A range of grating and crystal monochromators select various wavelength bands for some 20 different beam lines and associated user areas.

The calibration studies were carried out on the 3.4 VUV beam line which has a 30 – 1400 Å spectral range, with a computer controlled toroidal grating monochromator giving a 0.01 Å bandwidth. However this beam line is inefficient below 100 Å, and so while trends in response from 100 – 30 Å could be seen, an accurate calibration of the instrument over this range was impossible.

Light from the synchrotron ring selected by the monochromator was coupled from the high vacuum (10^{-9} torr) beam line to a lower vacuum (10^{-4} to 10^{-7} torr) experimental chamber via a 2 mm diameter glass capillary tube which acted as a pressure differential. The capillary tube was aligned by eye with the monochromator set to pass zero order light which has a strong visible component in the blue.

3.6.1 GRATING CALIBRATION

Figure 3.7 shows the experimental arrangement for the grating calibration. With the grating driven out of the synchrotron beam a vacuum photodiode biased to 18 V and connected to a high impedance (10^{14} Ω) pico-ammeter was used to characterise the photon flux as a function of wavelength. For a particular fill of the ring, flux is proportional to the electron beam current which slowly decays with a half life of 12 – 24 hours. Figure 3.8 is a typical trace of diode current versus wavelength showing the rapid cut off below 100 Å. The null current is the leakage current through the detector with the diode unilluminated.

The grating was then driven into the beam, and the blue component of the monochromator zero order used to set the grating angle to 3° grazing incidence. The vacuum photodiode was then used to measure the intensity of the various diffraction orders for a range of wavelengths. This was repeated with the grating rotated 90° about the beam axis to take account of the synchrotron light being highly polarised in the plane of the ring, and the two measurements were then averaged.

The ring current was noted before and after a series of measurements, and a correction made for the decay of the beam during the experiment. This was a small effect as the experiment typically took 10 – 20 minutes compared to the > 12 hour half life of the ring current. Figure 3.9 shows the measured grating efficiency as a function of wavelength for a range of diffraction orders.

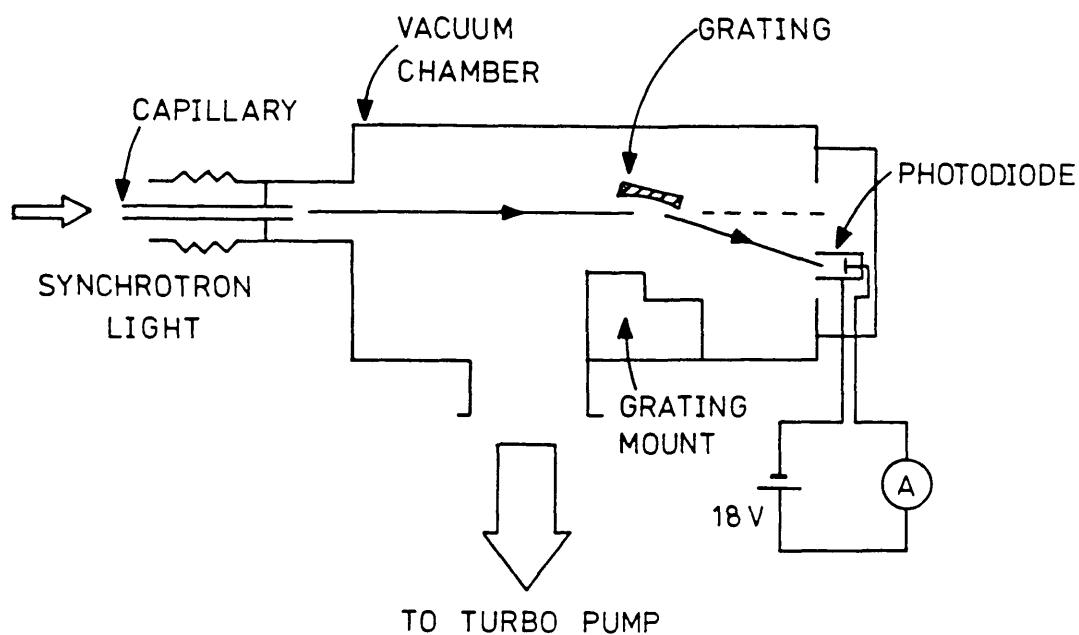


Figure 3.7 Experimental arrangement for the flat field grating calibration.

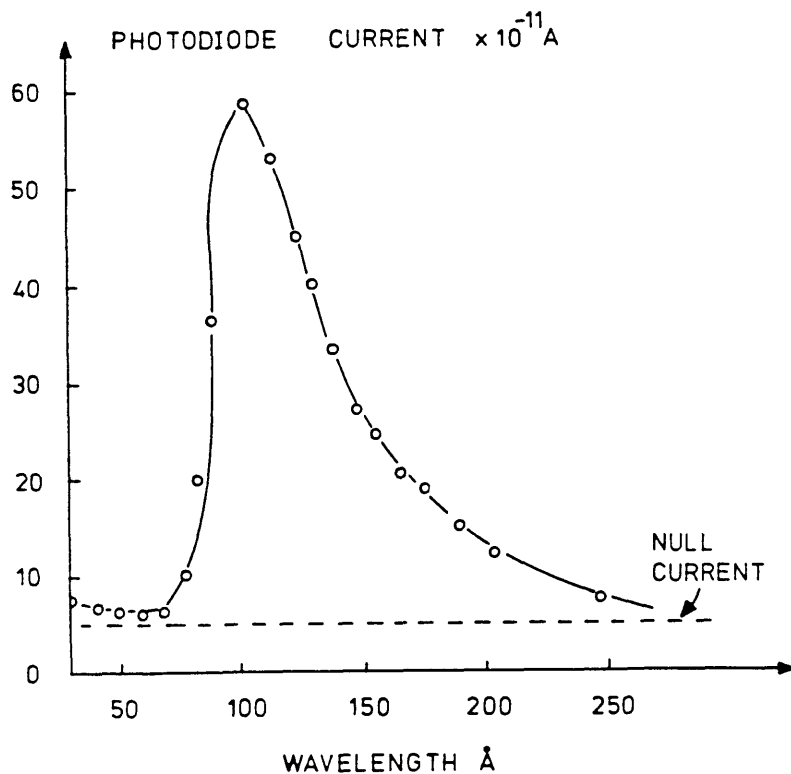


Figure 3.8 Vacuum photodiode response against wavelength selected by the VUV beam line monochromator.

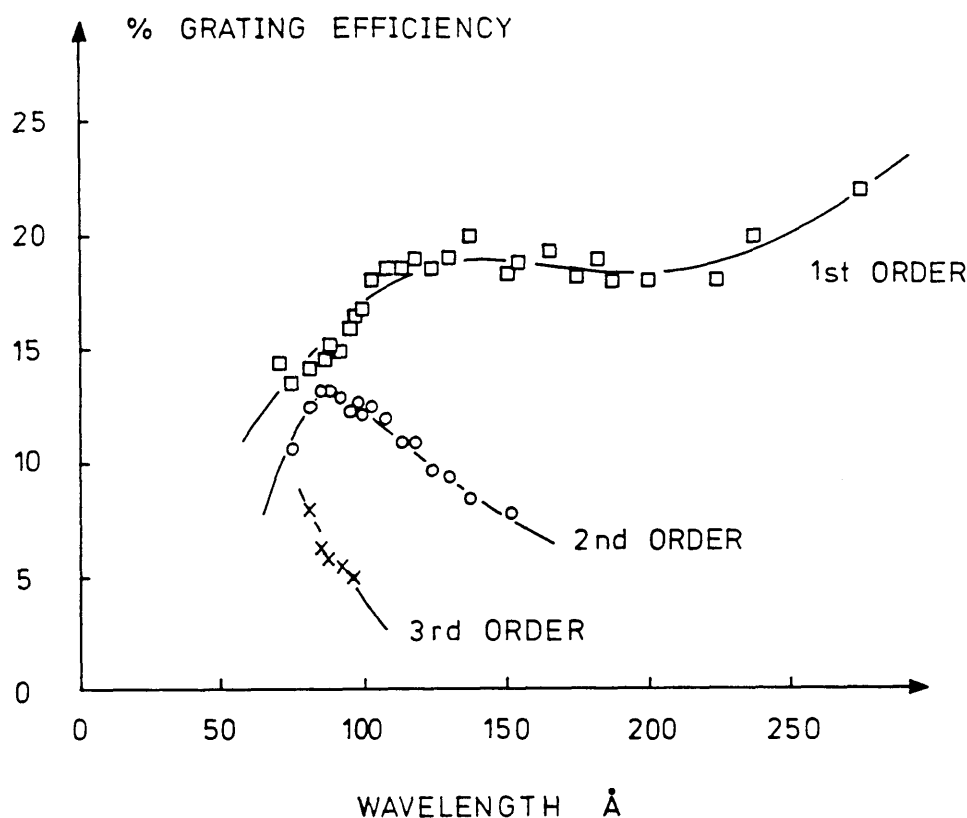


Figure 3.9 Measured grating efficiency versus wavelength for several diffraction orders.

The first order efficiency of the grating was well characterised, but the higher orders were not determined for the full wavelength range using this method. At short wavelengths the $n = 2$ and higher orders were too close together to distinguish with the photodiode, while at longer wavelengths only the first order was within the field of view of the photodiode. Figure 3.10 shows the $n = 1, 2, 3$ and 4 diffraction orders from a 1200 lines / mm grating at 71 Å, recorded with an XUV phosphor plate and image intensifier. We have found that at short wavelengths the grating efficiency does not fall off rapidly with diffraction order, contrary to the original characterisation carried out by the manufacturer using laser produced carbon plasmas [13]. This grating has been found to have a large high order efficiency, particularly at shorter wavelengths when used in conjunction with an XUV phosphor plate detector. For example 9th order fluorine He_α ($\lambda = 16.81$ Å) has clearly been observed in time integrated spectra (chapter 5, figure 5.4). This has caused significant problems with spectral identifications on some experiments, but has also been useful for separating lines in higher order which are beyond the resolving power of the instrument in 1st order (see for example chapter 6).

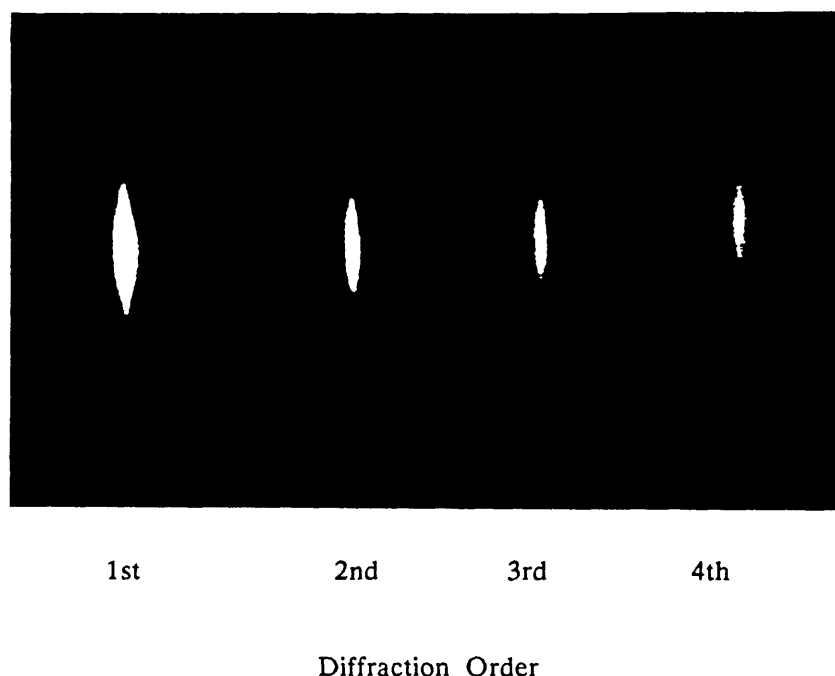


Figure 3.10 Picture of the $n = 1, 2, 3$, and 4 diffraction orders recorded at $71 \text{ \AA}$

3.6.2 FILTER CALIBRATION

A range of filters are used on laser plasma experiments, both to protect the grating from plasma blow off and debris damage, and to differentially filter the spectrum and enhance a particular wavelength band. A series of formvar (polyvinyl formal resin $[-\text{CH}_2\text{CH}(\text{OOCH}_3)-]_n$) and aluminium foil filters were characterised over a range of wavelengths using a similar method to the grating calibration.

A series of readings of synchrotron brightness with and without filters in front of the photodiode were taken for several wavelengths. Again a small correction for the decay of the ring current during the experiment was made. Figures 3.11, 3.12 and 3.13 show the measured and calculated transmission curves for $1000 \text{ \AA}$ formvar, $4000 \text{ \AA}$ formvar and $8000 \text{ \AA}$ aluminium filters as a function of wavelength. The $1000 \text{ \AA}$ formvar filter shows a particularly good fit to the theoretical transmission curve, demonstrating the precision with which the Target Preparation Group at RAL make these filters.

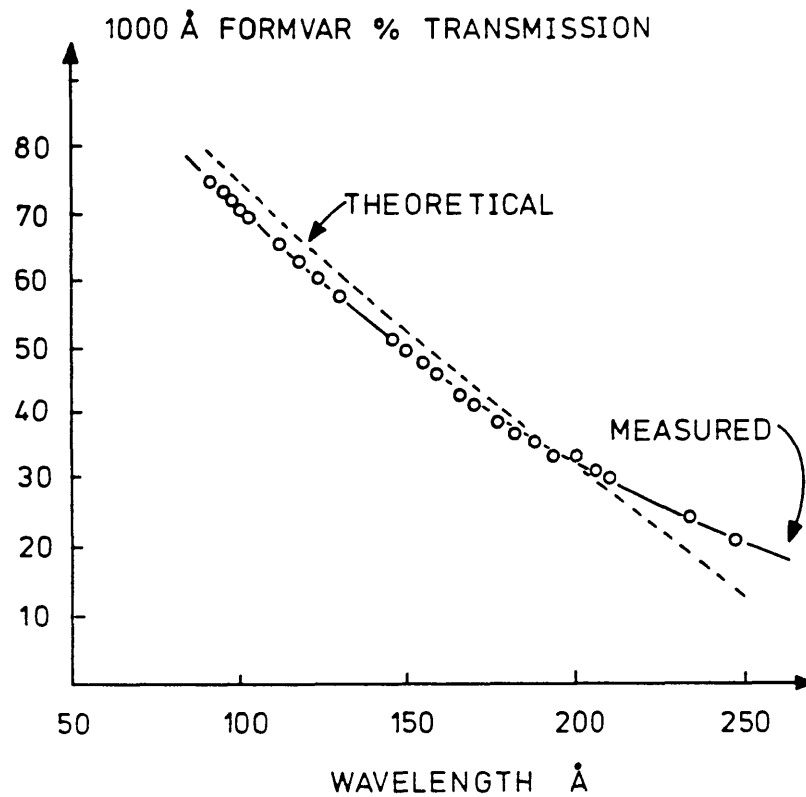


Figure 3.11 Calculated and measured transmission curves as a function of wavelength for a 1000 Å Formvar filter.

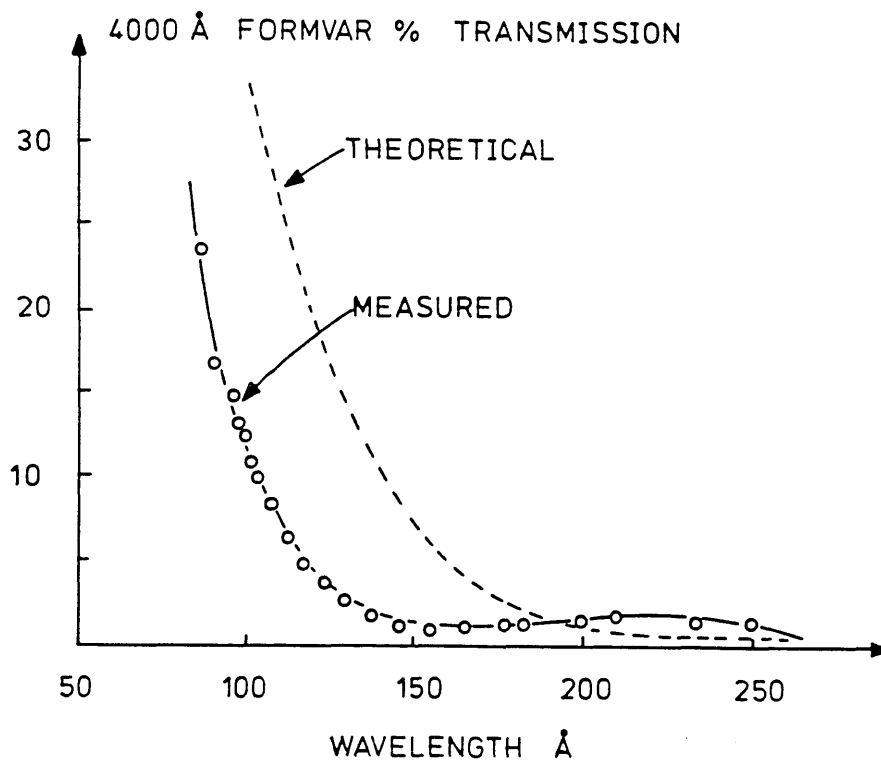


Figure 3.12 Calculated and measured transmission curves as a function of wavelength for a 4000 Å Formvar filter.

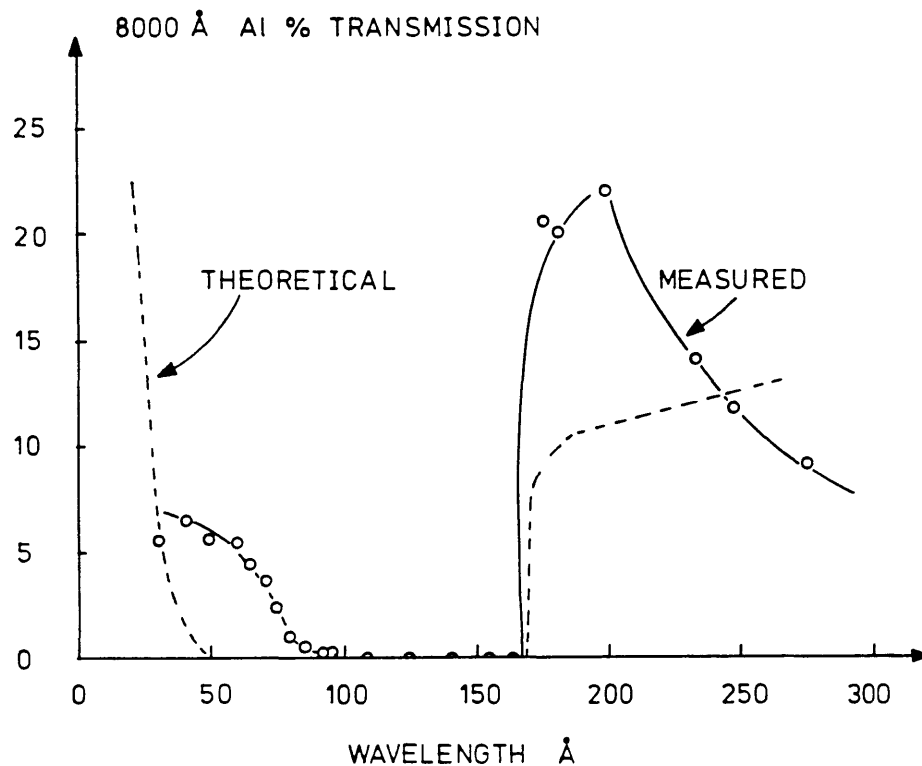


Figure 3.13 Calculated and measured transmission curves as a function of wavelength for an 8000 Å aluminium filter.

3.6.3 PHOSPHOR PLATE CHARACTERISATION

XUV phosphor plates coupled to image intensifiers are used to record time integrated spectra from the flat field spectrometer. A phosphor plate / image intensifier pair was characterised by allowing the beam from the monochromator (previously characterised by the photodiode) to hit the phosphor plate directly. For a range of wavelengths HP5 film was used to record the output from the intensifier triggered in single pulse mode. A density step wedge was printed onto film and was used to deconvolute the film response after developing in D19 at 23°C for 4 minutes.

Figure 3.14 shows a typical microdensitometer trace of the density step wedge and the film density versus log (exposure) curve for HP5 film, while figure 3.15 shows the number of XUV photons required to produce a unit of specular film density ($D = 0.07$) as a function of wavelength for the phosphor plate and intensifier pair.

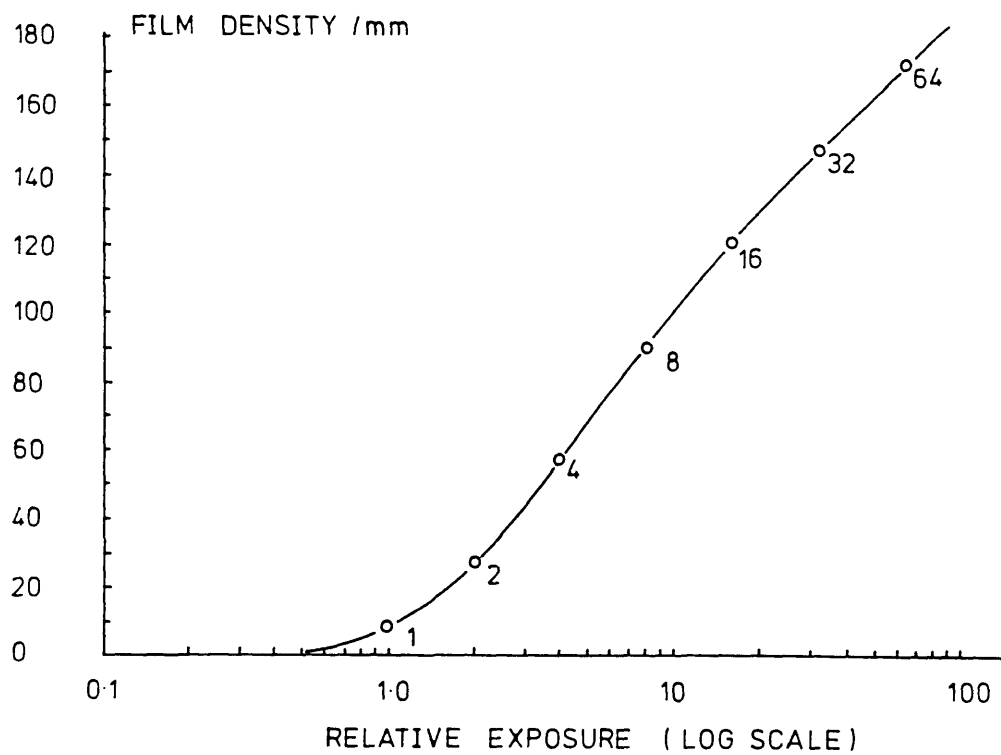


Figure 3.14 Log / linear plot of film density versus exposure for a calibration step wedge.

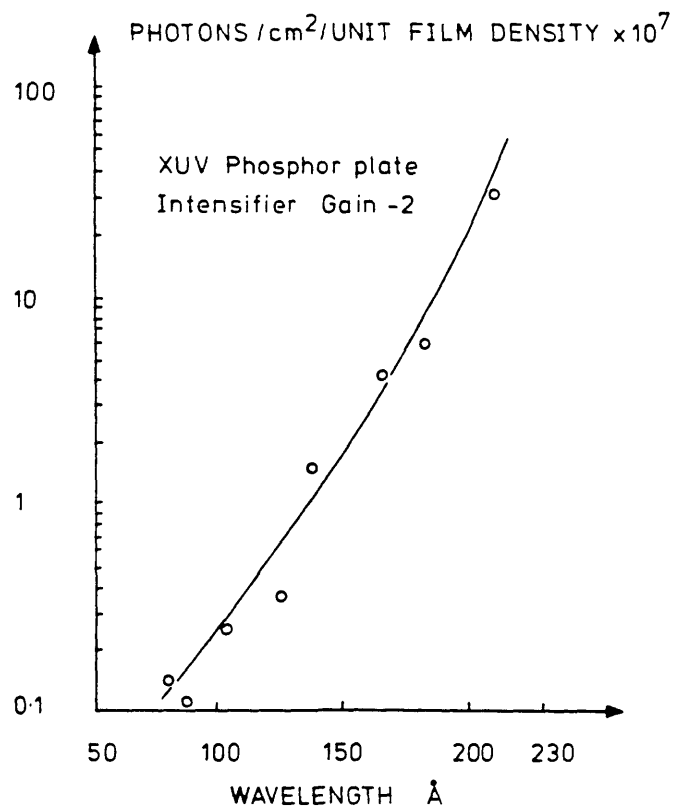


Figure 3.15 Absolute calibration of the phosphor plate as a function of wavelength.

As can be seen from figure 3.15 the number of XUV photons needed to produce a given density on film falls quite rapidly for shorter wavelengths. This is due in part to an aluminium coating on the phosphor plate used to block visible light, and results in an enhancement of short wavelength high order lines in time integrated spectra (see chapter 4, figure 4.6 for example).

3.6.4 STREAK CAMERA AND PHOTOCATHODE CALIBRATION

The streak camera and image intensifier unit used to time resolve flat field spectra was calibrated in single pulse mode in much the same way as the phosphor plate. A beam from the synchrotron monochromator hit the streak camera photocathode directly, and the output from the streak camera was recorded on HP5 film via an image intensifier. A range of high and low density CsI photocathodes were characterised in this way, and the most sensitive found to be 10 μm of low density CsI coated onto a 50 $\text{\AA}$ aluminium, 1000 $\text{\AA}$ formvar, 5 μm copper wire mesh substrate.

Figure 3.16 shows the number of photons required to give unit specular density ($D = 0.07$) on HP5 film as a function of wavelength. The shape of the response curve is due primarily to the efficiency of CsI as a photoelectron emitter for different wavelengths. This has been examined in some detail by groups working on streak camera design [72,73]

3.6.5 TIME RESOLVED SPECTROMETER CALIBRATION

The complete time resolved spectrometer unit was characterised by convoluting the wavelength response of the various individual components. Figure 3.17 shows the number of photons required to give unit density ($D = 0.07$) on HP5 film as a function of wavelength for the 1000 $\text{\AA}$ formvar filter, 1200 lines / mm grating, 10 μm low density CsI photocathode, 135 ps / mm sweep speed configuration typically used on an X-ray laser experiment.

The response curve is thought to be accurate to within a factor of 2, the primary sources of error being due to the difficulty in measuring photodiode currents on the order of a few times 10^{-11} A in the electrically noisy environment of the synchrotron, and the aging of the NBS reference diode. This calibration has since been used to determine the absolute power output of an experimental 182 $\text{\AA}$ hydrogenic carbon laser (chapter 4).

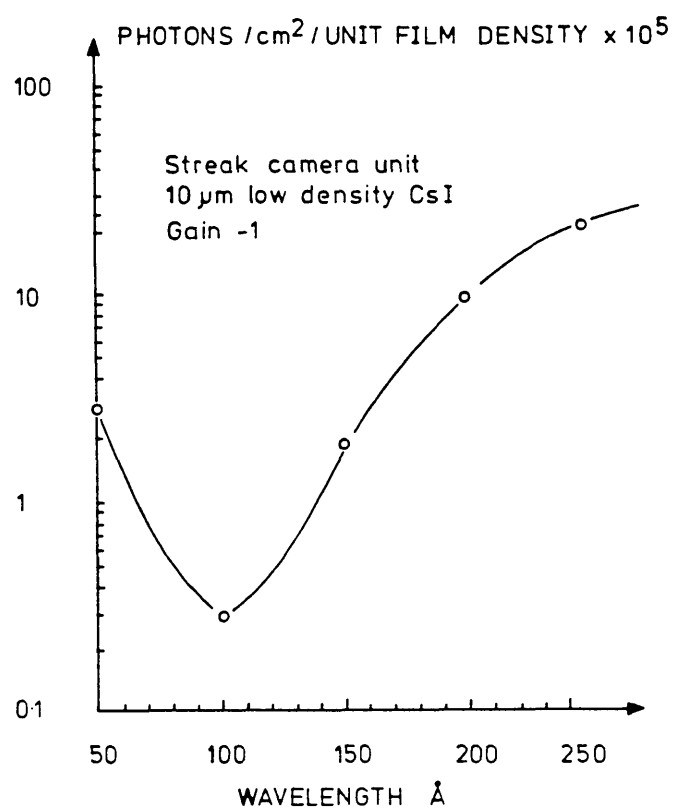


Figure 3.16 Absolute calibration of the CsI photocathode, streak camera and image intensifier unit.

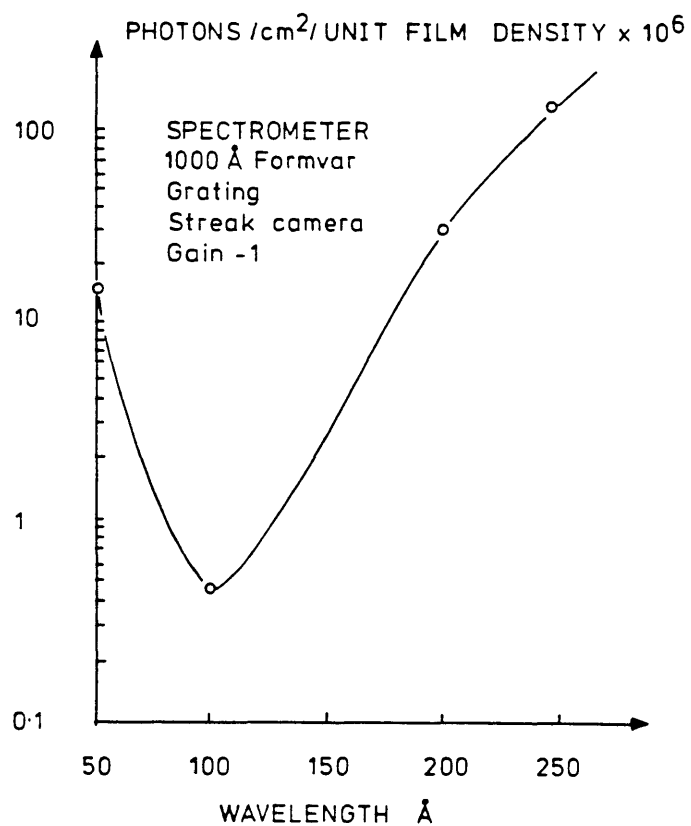


Figure 3.17 Absolute calibration of the complete time resolving XUV spectrometer as a function of wavelength.

CHAPTER 4. HYDROGENIC CARBON X-RAY LASER EXPERIMENTS

ABSTRACT

The fibre based hydrogenic recombination XUV laser scheme was investigated experimentally [8]. Carbon fibres 7 μm in diameter and up to 9 mm in length were irradiated in line focus geometry using the VULCAN glass laser at RAL. Time resolved XUV spectroscopy of the axial emission from fibre targets was used to demonstrate exponential growth of intensity of the carbon Balmer H_α line at 182 $\text{\AA}$ with plasma length. A gain coefficient of $3.35 \pm 0.33 \text{ cm}^{-1}$ at coupled energies of $\approx 2.6 \pm 0.6 \text{ J cm}^{-1}$ was inferred from axial streak spectra of the freely expanding adiabatically cooled plasmas.

A second experiment carried out by the author used simultaneous axial and transverse spectroscopy with improved temporal resolution to demonstrate gains of up to 6.1 cm^{-1} on a single shot. Discrepancies between theoretical predictions for the time evolution of gain and experimental results from axial measurements alone were also resolved with this technique. Good fits were found between improved numerical simulations and the gain given by axial to transverse ratios.

4.1 INTRODUCTION

The adiabatically cooled hydrogenic carbon recombination XUV laser scheme outlined in chapter 2 has been investigated experimentally by a number of previous workers. Population inversions were observed in laser produced carbon plasmas by Irons and Peacock [32] and also by Perts group [33], but at densities too low to give significant gain. Pert and co-workers have reported gains of $\approx 2 - 4 \text{ cm}^{-1}$ in plasmas from laser irradiated fibres [34] but over relatively small lengths ($< 2 \text{ mm}$) using time integrated diagnostics. The experiments described in this chapter include significant advances in both diagnostic techniques and plasma lengths compared to previous investigations.

The conditions for maximum gain identified by the semi-analytical arguments presented in chapter 2 were refined using detailed numerical simulations [74,75]. Equation (2.51) identifies small fibre diameters and short pump laser pulses for maximum gain. However the fibre diameter is restricted to $\geq 7 \text{ }\mu\text{m}$ as smaller fibres become too unstable to align well. Similarly the minimum pulse length of $\approx 70 \text{ ps}$ is constrained by the range of the VULCAN laser short pulse oscillator. With these experimental limitations ($r_0 = 3.5 \text{ }\mu\text{m}$, $t = 70 \text{ ps}$) and a pump laser wavelength $\lambda = 0.53 \text{ }\mu\text{m}$, equations (2.50) and (2.51) identify an absorbed energy of 1.6 J cm^{-1} for peak gain. Simulations refined this to 2.5 J cm^{-1} which for coupling efficiencies of $\approx 10 \%$ into fibre targets with the TAE line focus gives a pump laser requirement of 25 J cm^{-1} . This can be achieved using only one pair of opposing beams in the six beam TAE line focus, which allows lengths up to 21 mm to be irradiated at the required level by spacing the three pairs of beams along the target.

4.2 EXPERIMENTAL (AXIAL SPECTROSCOPY ONLY)

Carbon fibres $7 \text{ }\mu\text{m}$ in diameter and supported at one end were irradiated with 70 ps pulses of green ($0.53 \text{ }\mu\text{m}$) light in line focus geometry. Targets were aligned under vacuum to $\pm 5 \text{ }\mu\text{m}$ in position and 10^{-3} in angle using a split field microscope system [68] and motorised target mount. Lengths greater than $\approx 12 \text{ mm}$ were found to be subject to vibration and gravitational bending which made this mounting arrangement unsuitable for longer fibres. Various mechanical mounting techniques have been tried for longer fibres, none of these has been particularly successful.

It has also been suggested [76] that the electrostatic force from a 3–4 mm diameter ring charged to a few hundred Volts, placed beyond the free end of the fibre would both support the fibre and damp out oscillations. However this technique has yet to be used on an experiment.

Four beams of the line focus were then individually aligned on target using a CW argon ion laser, by adjusting the line focus mirror to minimise the light transmitted past the fibre. This was carried out under vacuum to avoid refraction errors. For lengths ≤ 7 mm two opposing beams were used, and some of each beam masked off. For lengths between 7 and 14 mm two pairs of opposing beams were used. Care was taken to ensure that the free end of the fibre which points towards the XUV spectrometer was fully illuminated, as cold plasma at this point would reabsorb any ASE signal from the bulk of the plasma.

The primary diagnostic was a time resolved flat field 1200 lines / mm XUV spectrometer [14] aligned along the axis of the fibre target. The spectrometer observed line emissions from the plasma with $0.5 \text{ \AA}$ resolution over a $100 \text{ \AA}$ spectral window, with ≈ 380 ps temporal resolution over a 3.75 ns time interval. The time resolution of the instrument was set by the product of the streak camera sweep rate, input slit width and magnification ($135 \text{ ps / mm} \times 2 \text{ mm} \times 1.4 \approx 380 \text{ ps}$). The spectrometer was run in the slitless mode to maximise the acceptance angle, and was protected from plasma blow off by a $1000 \text{ \AA}$ formvar filter.

A range of additional diagnostics were used to characterise the laser produced plasmas (see figures 4.1 and 4.2). The plasma length was recorded with an active pinhole camera comprising an $80 \text{ }\mu\text{m}$ platinum pinhole, phosphor plate and image intensifier. The camera was filtered with $0.8 \text{ }\mu\text{m}$ of Al to pass $\geq 620 \text{ eV}$ ($\leq 20 \text{ \AA}$) Bremsstrahlung photons emitted during the laser heating phase (the shortest carbon line is 1s–6p at $26 \text{ \AA}$). A second unfiltered passive pinhole camera recording directly onto DEF X-ray film was used as a back-up.

The plasma uniformity along the length of the target was investigated using a 1D space resolving, time integrated temperature monitor. This instrument used a series of 15, 30 and $45 \text{ }\mu\text{m}$ thick beryllium filters to remove carbon line emission ($\leq 10^{-4}$ transmission) and recorded the continuum emission for photon energies of $h\nu > 1200 \text{ eV}$, $h\nu > 1450 \text{ eV}$ and $h\nu > 1650 \text{ eV}$ with a phosphor plate and image intensifier. 1D space resolution was provided by a 1 mm slit placed in front of the filter-pack.

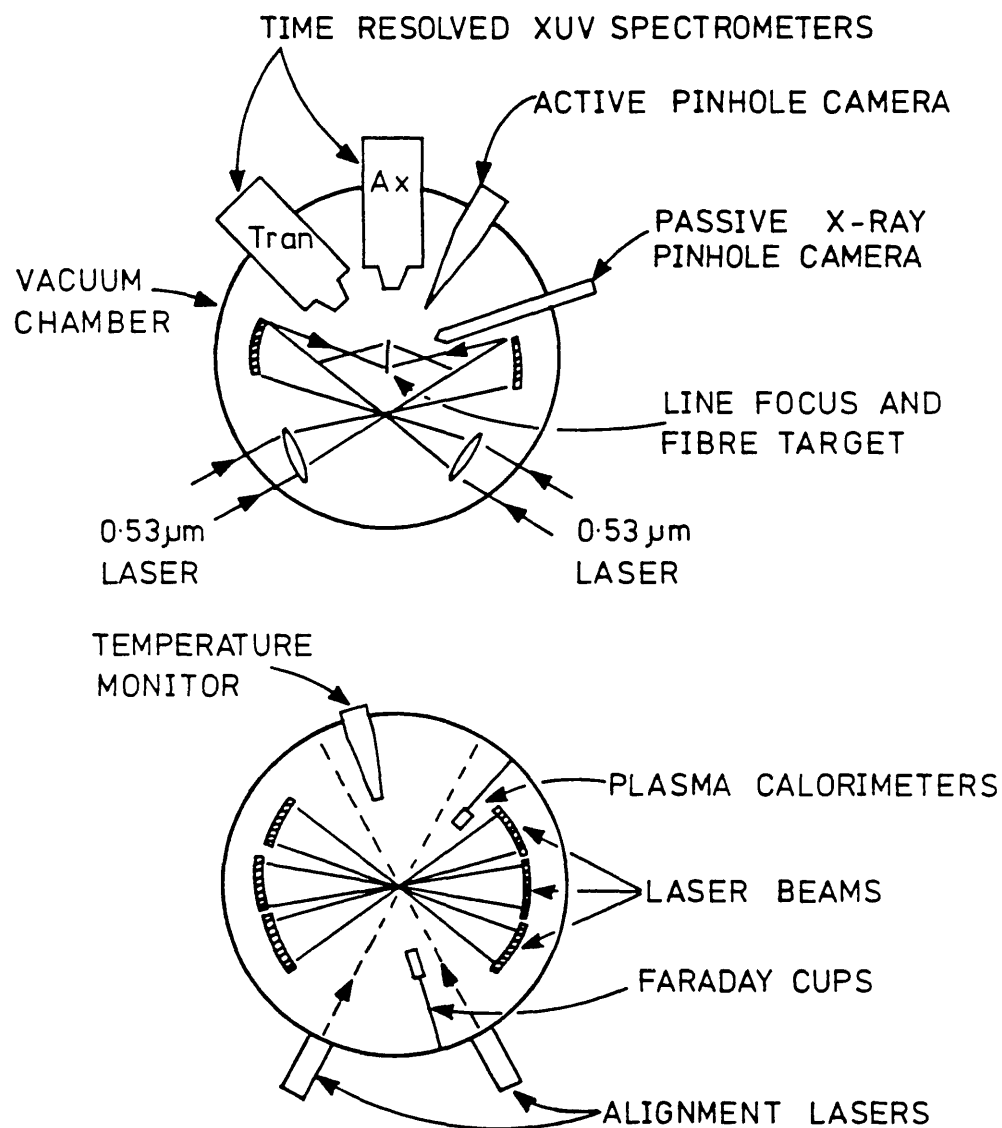


Figure 4.1 Experimental arrangement for the carbon fibre X-ray laser studies.

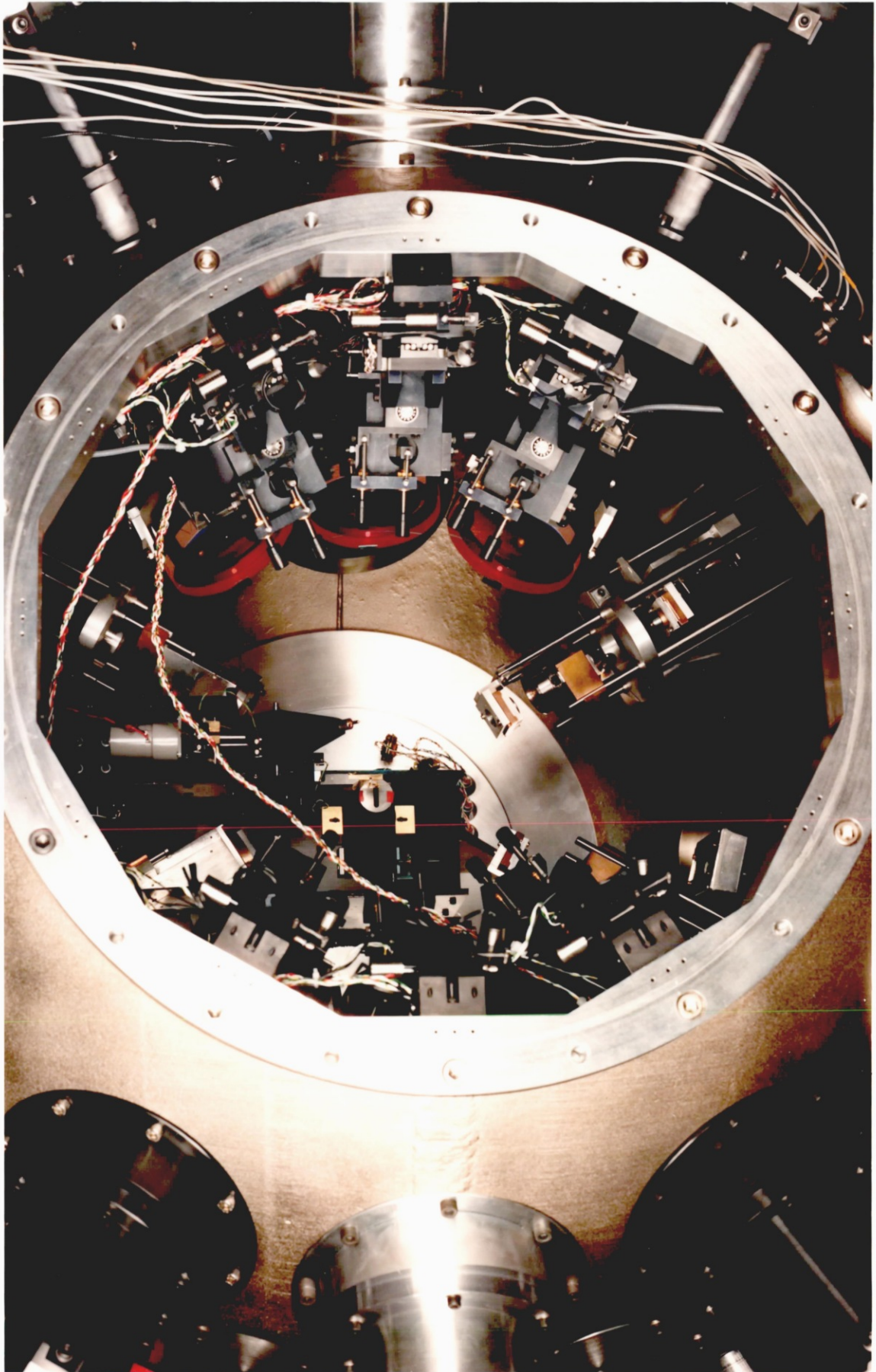


Figure 4.2 Picture of the diagnostics and target chamber for X-ray laser experiments.

The temperature at any point in the image could be found from the continuum slope. While this diagnostic was time integrated, the relatively hard X-rays which it records are only emitted at early times when the plasma is hot and dense, and so it is a good indicator of the plasma uniformity immediately after the heating pulse.

The absorbed energy was measured with an array of plasma calorimeters. Each element of the array consisted of two 3x3 mm by 10 μm thick Al foils spot welded with 25 μm diameter Constantan / Chromel thermocouple wires. The two foils were connected to a differential millivoltmeter, and one foil in each pair was screened from plasma blow off with a transparent cover. Scattered pump laser light was seen equally by both foils, and was subtracted out of the signal by the differential millivoltmeter. The plasma however hit only the unscreened foil giving a net response from the element.

4.3 PLASMA CHARACTERISTICS

The line focus arrangement gave good uniformity of illumination along the full length of the fibre targets, even though beam pairs were not overlapped. This can be seen in figure 4.3 which shows a typical active pinhole camera (APC) image from a 4.5 mm long plasma. The plasma length was taken to be that from which 90 % of the X-rays in the APC image were emitted. This method ignores the relatively cold and tenuous plasma which expands away from the ends of the fibre along the axis, and which does not give any significant contribution to the gain of the system.

Figure 4.4 shows the distribution of plasma energy parallel to the fibre axis recorded by the plasma calorimeter array on a 7.2 mm shot. The total plasma energy was found by integrating under the curve and assuming that the plasma was radially symmetrical. The assumption of radial symmetry is not particularly valid in this regime as the fibres are only irradiated from the sides. The fibres are too large for lateral heat transport to heat the unilluminated top and bottom of the fibres to the same degree as the sides, and the plasma expands outwards as two lobes rather than a cylinder. However this assumption is believed to result in an error within the absolute accuracy of the calorimeters. The average absorbed energy per unit length was found to be $2.6 \pm 0.6 \text{ J cm}^{-1}$.



Figure 4.3 Active pinhole camera image of a 4.5 mm long carbon plasma showing good plasma uniformity.

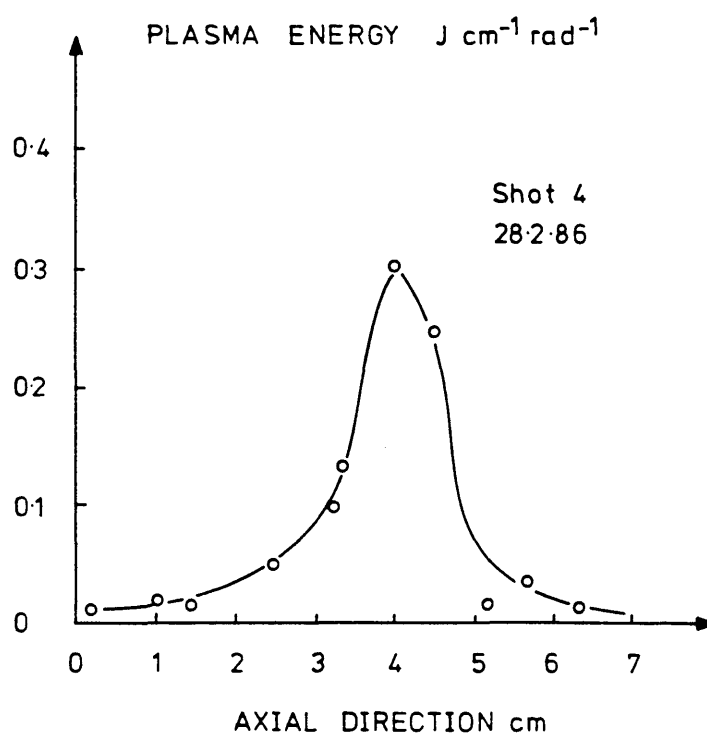


Figure 4.4 Plasma energy distribution parallel to the fibre axis on a 7.2 mm carbon fibre shot.

The average time integrated electron temperature found from the continuum slopes recorded by the 1D imaging temperature monitor was 250 ± 100 eV in the hotter regions of the plasma. Some degree of variation along the fibre was visible in the temperature monitor images (see figure 4.5) and this was thought to be due to both irregularities on the fibre surface and non-uniformities in the pump laser beams which are not of a particularly high quality. However the large errors in the results from this diagnostic limit its usefulness as an absolute measure of the electron temperature at the start of the expansion phase.

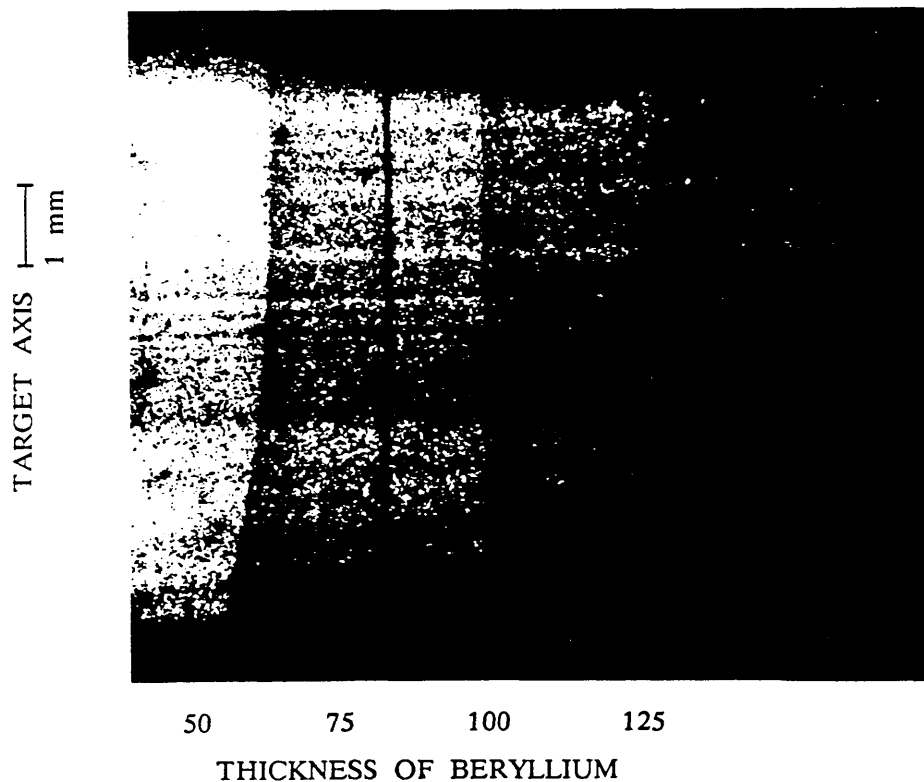


Figure 4.5 Space resolving temperature monitor image showing some evidence of variation along the fibre axis

4.4 AXIAL XUV SPECTROSCOPY

The axial spectrometer was initially run up in the time integrated mode with a phosphor plate detector. The phosphor plate covers the full spectral range of the grating and this allows for easier identification of the spectrum. It is also more convenient to check the alignment of the spectrometer in this mode before timing up the streak camera. Figure 4.6 shows a typical time integrated spectrum from a CH slab target for which no gain is expected. The spectrum is dominated by higher order H like carbon resonance lines, and this is due to the higher sensitivity of the phosphor plate at shorter wavelengths (see section 3.6.3). The presence of such strong hydrogenic resonance lines indicates that the carbon plasma was fully ionised during the short heating pulse. The laser line carbon Balmer H_{α} at 182 Å is visible, but is one of the weakest lines in the spectrum.

A 100 Å section of the spectrum was time resolved by replacing the phosphor plate with a Kentech high magnification streak camera [70] incorporating a low density CsI XUV photocathode. Figure 4.7 shows a streak spectrum and densitometer trace from a 5.6 mm long carbon fibre target. The spectrum is now dominated by the Balmer series with H_{α} , H_{β} , H_{γ} and H_{δ} being the most prominent features. The weaker lines are H and He like resonance transitions in higher order. The short burst of continuum emission at time $t = 0$ is due to Bremsstrahlung from the plasma during the laser heating phase. Line emission turns on nearly coincident with the Bremsstrahlung but does not peak until several hundred picoseconds after the heating pulse when the plasma has begun to expand and has cooled significantly, enhancing the recombination.

The brightest line in the spectrum appears to be carbon H_{β} at 135 Å rather than the expected laser line H_{α} at 182 Å. No significant gain is expected for H_{β} even though the $n = 4-2$ levels are likely to be inverted, since the oscillator strength for this transition is too low. This result is due simply to the wavelength response of the spectrometer and streak camera. From figure 3.17 which gives the wavelength calibration for the whole instrument it can be seen that the differential sensitivity from 135 Å to 182 Å is $\approx 10 : 1$.

TIME INTEGRATED CARBON SPECTRUM

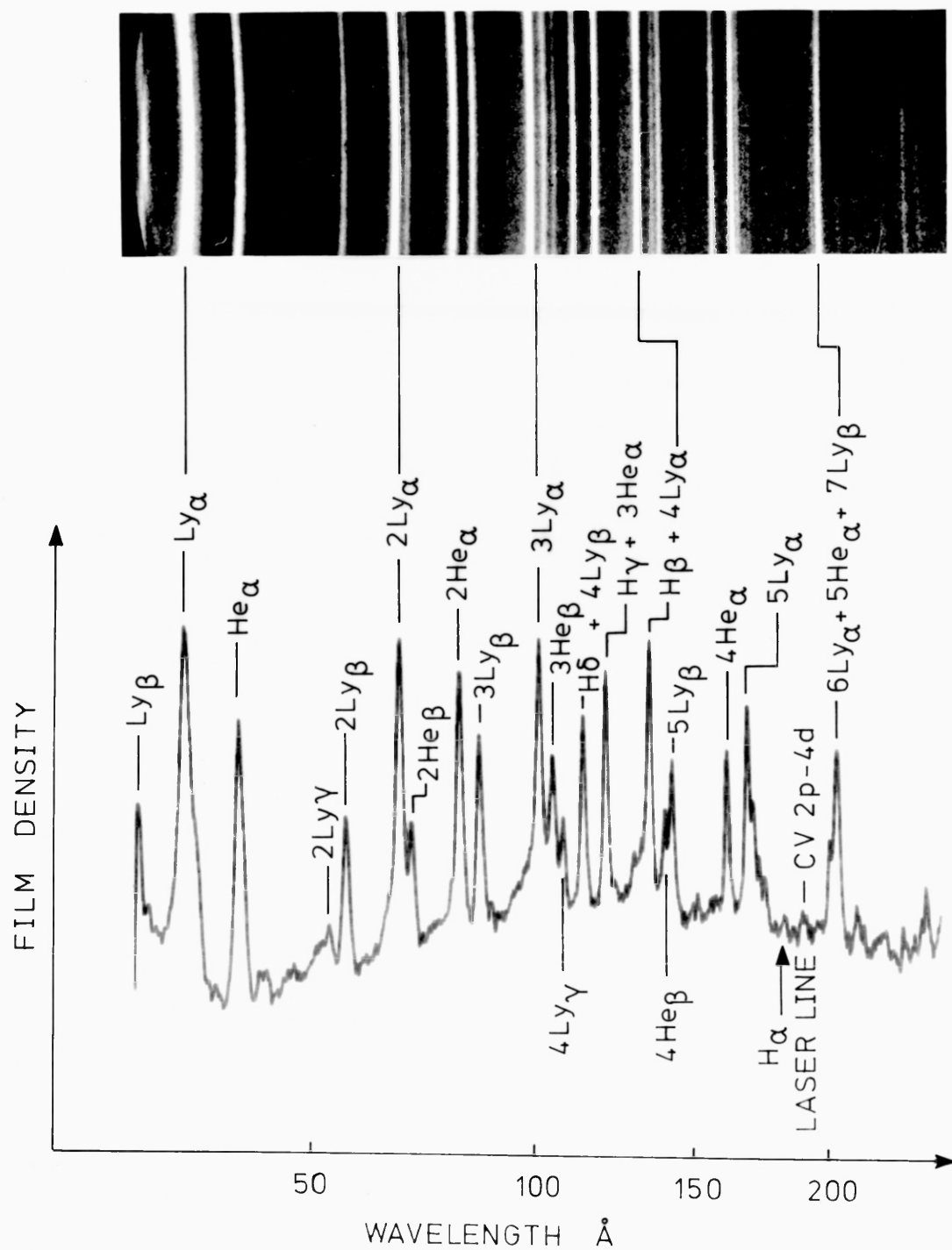


Figure 4.6 Typical time integrated XUV spectrum from a CH slab target.

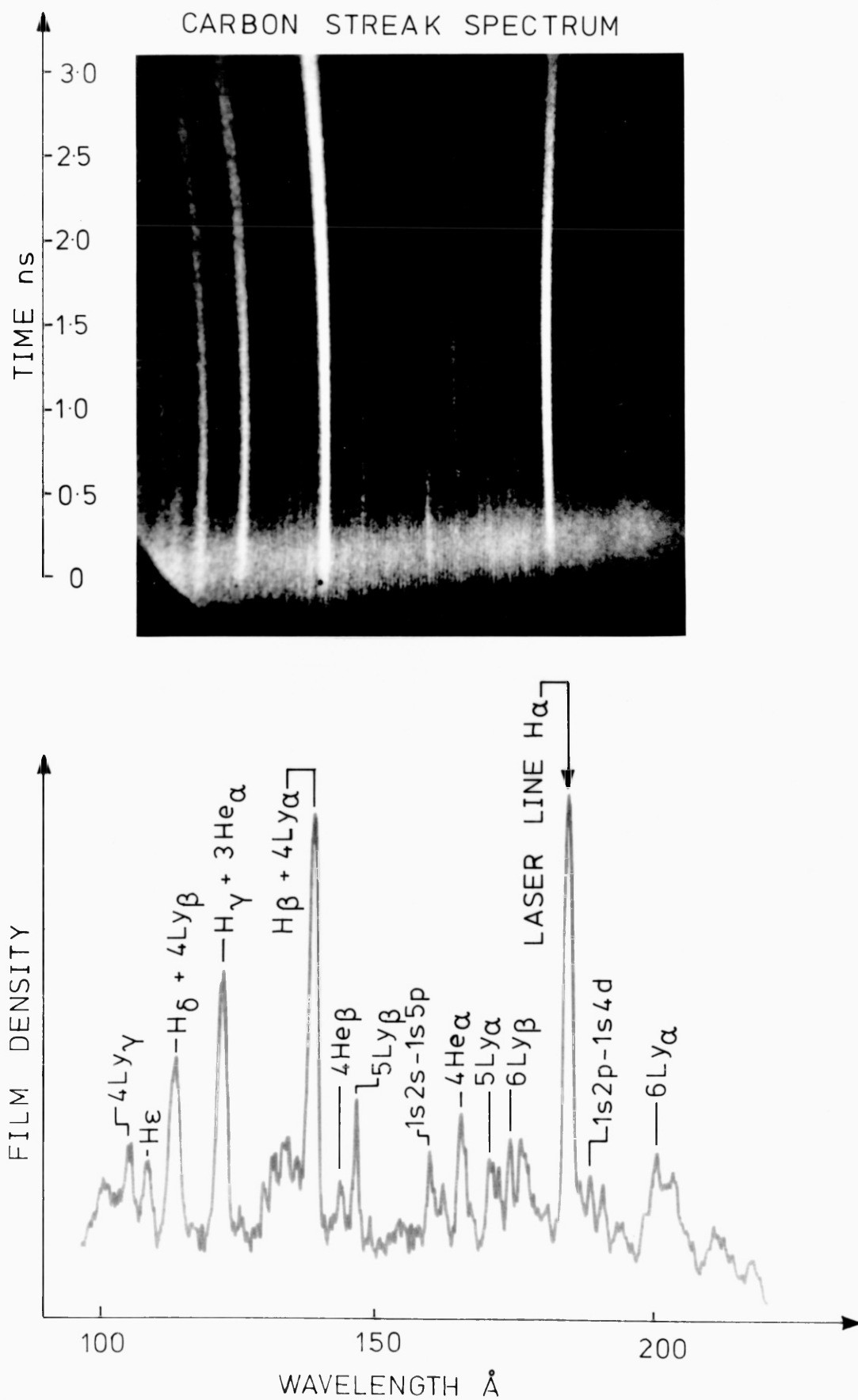


Figure 4.7 Axial XUV streak spectrum and densitometer trace from a 5.6 mm long carbon fibre shot.

There are also higher order contributions to some of the Balmer series which changes their apparent ratios. This effect was investigated using a 2000 Å Al, 1000 Å formvar filter to suppress the Balmer series on some shots. The H_α laser line was found to be free from any higher order contributions. The line at 135 Å was composed of 71 % H_β + 29 % Ly_α , the line at 120 Å was 20 % H_γ + 80 % $3He_\alpha$ and the line at 113 Å was 24 % H_δ + 76 % $4Ly_\beta$. These results were confirmed using an 8000 Å Al foil filter to suppress all lines between 50 and 165 Å ($\leq 10^{-6}$ transmission) leaving only the resonance lines.

The H_α emission peaked between 0.8 and 1.2 nanoseconds after the start of the laser pulse and evidence for gain was looked for at this time in the axial streak spectra. Non linear growth of H_α with plasma length was immediately obvious from the axial spectra. Figure 4.8 shows streaks from 1.9 and 6.8 mm carbon fibre targets. In the densitometer traces of these spectra (figure 4.9) it can clearly be seen that the ratio of H_α to H_β on film changes significantly from 1 : 2 to 1 : 1 with increasing plasma length, and this corresponds to an absolute change of 5 : 1 to 10 : 1 when folding in the instrumental sensitivity and film curve.

The variation of H_α and H_β intensity with plasma length is shown for the full data set of some 20 shots in figure 4.10. Film and instrumental responses have been deconvoluted from this data to give absolute fluxes in photons / second steradian. The H_α transition shows strong exponential growth of the form $I \propto \exp(\alpha \ell) - 1$ with a gain coefficient of $\approx 4.78 \pm 0.37 \text{ cm}^{-1}$, while H_β for which no gain is expected shows a near linear increase with plasma length. The gain coefficient for H_α was found using a log / linear least squares fitting to the axial data. The error bars shown in figure 4.10 are due to the uncertainty in the plasma lengths measured by the active pinhole camera, and the errors in H_α intensity from the uncertainty in the level of the background emission which underlays the line emission. A number of points lie several σ away from the best fit gain curve and this is due to shot to shot variation in coupled energy. This effect is difficult to correct for using axial spectroscopy alone.

There is some small non linearity in both the H_β and H_γ intensity variation with plasma length. This is thought to be caused by a systematic error resulting from some of the side light emission from the fibre being seen by the axial spectrometer. Scaling the axial data to give a linear growth of H_γ reduces the observed H_α gain from 4.7 to $3.35 \pm 0.33 \text{ cm}^{-1}$ as can be seen in figure 4.11. This represents a lower limit for the experimentally measured gain.

AXIAL STREAK SPECTRA

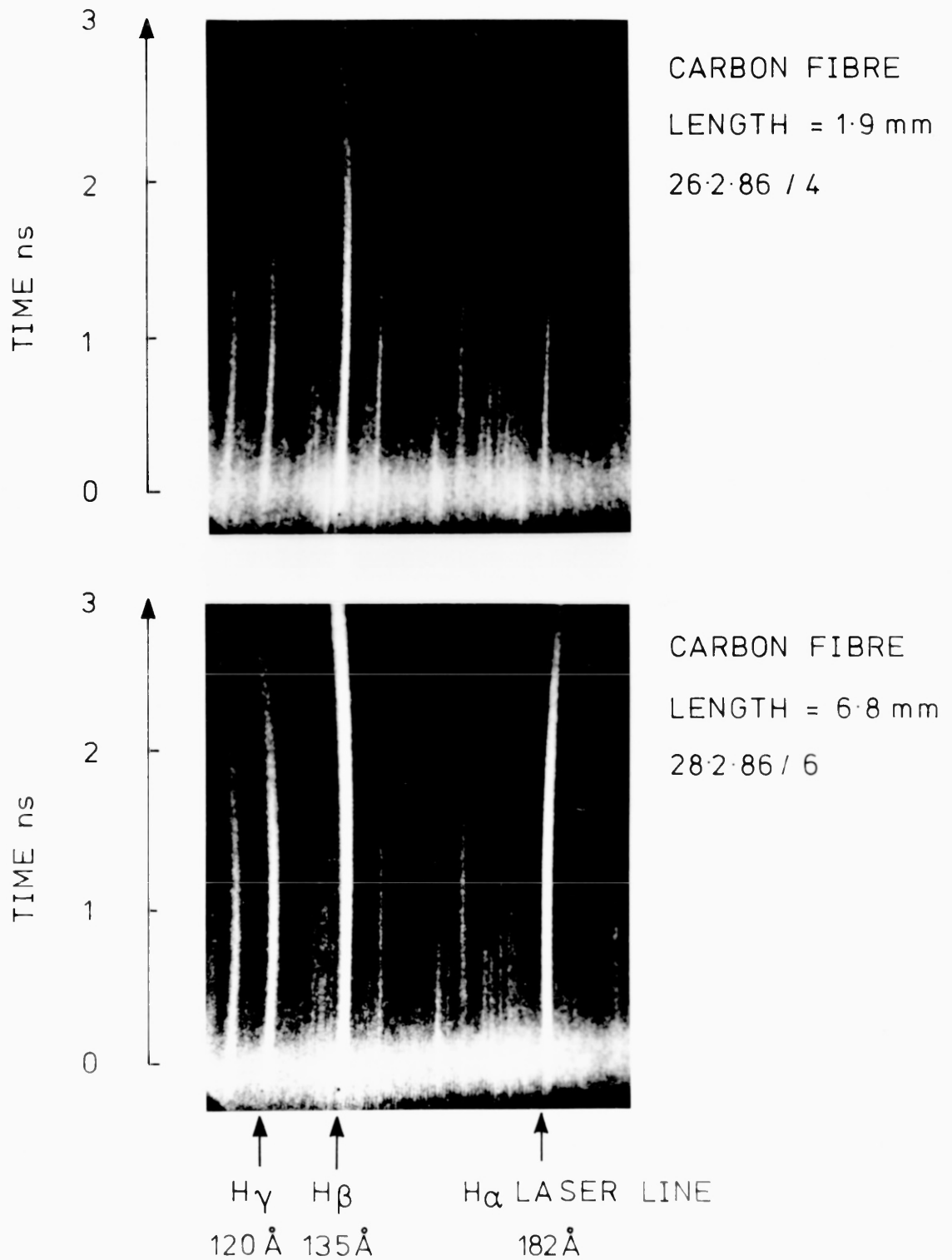


Figure 4.8 Axial streak spectra from 1.9 and 6.8 mm carbon fibres

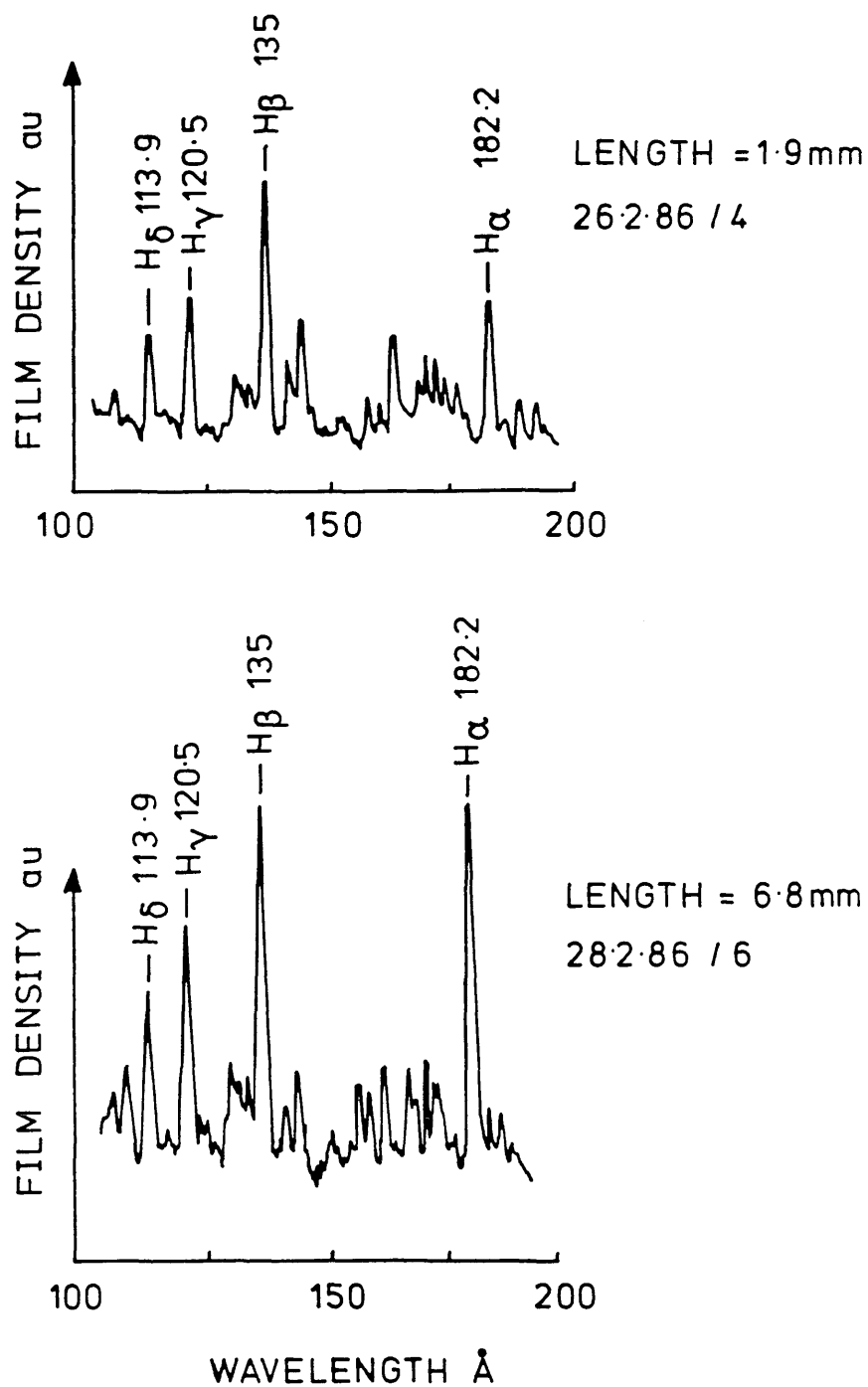


Figure 4.9 Densitometer traces of the streak spectra presented in figure 4.8

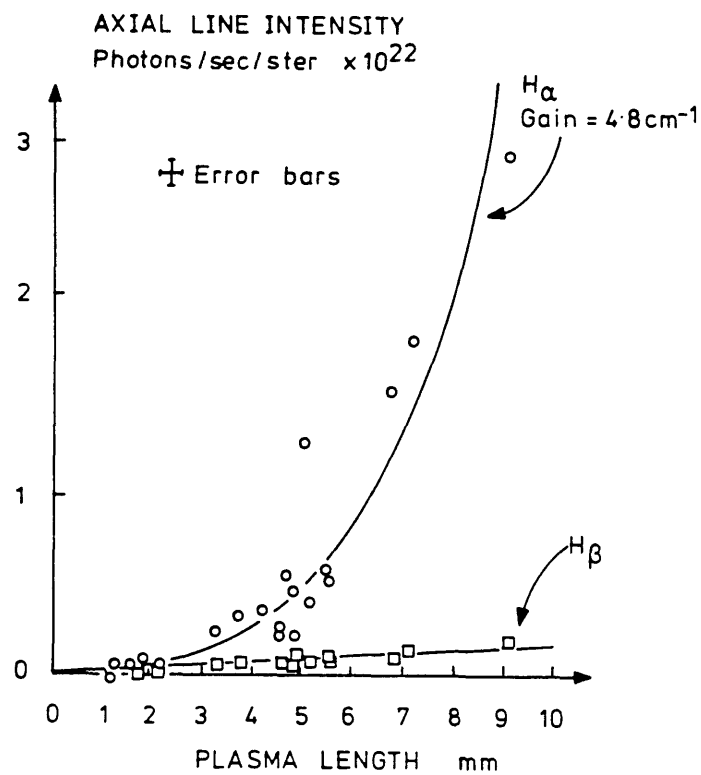


Figure 4.10 Variation of axial H_α and H_β intensity with plasma length for the full data set.

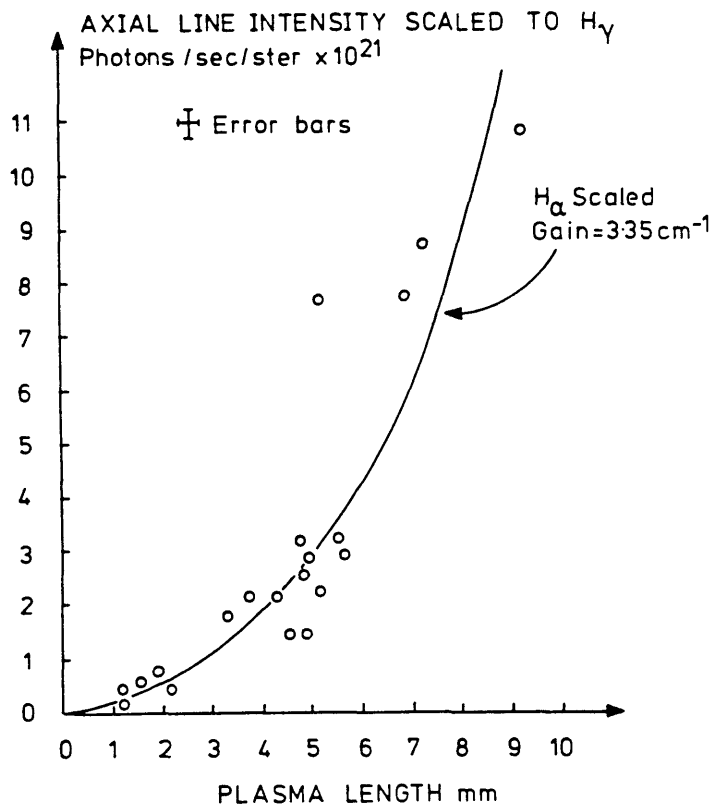


Figure 4.11 Axial H_α intensity versus plasma length, with the data scaled to give a linear growth of H_γ .

The power output of the carbon ASE "laser" can be estimated using the absolute calibration of the spectrometer and streak camera described in section 3.6.5. For the longest recorded shot of 9.1 mm (8.3.86 No. 12) this gives an approximate power output of 380 KW / steradian. Integrating the output over time gives a total energy of ≈ 1 mJ and from this we can estimate the efficiency of the laser to be $\approx 4 \times 10^{-3} \%$ by assuming a pump laser input energy of 25 J. All these results are only accurate to within a factor of 2 because of the uncertainty in the spectrometer calibration.

It is important to note that the exponential growth of H_α is not obvious for lengths less than 6 mm. This places an experimental limitation on the short wavelength scalability of this scheme when using the VULCAN laser as a pump. Lengths greater than 6 mm should be investigated before clear conclusions about the existence of gain can be drawn. However the $I \propto Z^4$ scaling of pump laser irradiance which accompanies the Z^{-2} scaling of XUV laser wavelength will rapidly use up all our available pump power.

4.5 TIME EVOLUTION OF GAIN FROM AXIAL SPECTRA

The time evolution of the Balmer series can also provide a good indication of the existence of gain. Figure 4.12 compares the time evolution of H_α and H_β from 1.8 and 4.9 mm shots. In the short shot H_α and H_β show similar temporal behaviour as the product $\alpha\ell$ is small. However in the longer shot H_α is significantly peaked at early times and while this is not conclusive evidence for gain, it does support the exponential length variation seen in axial spectra.

The time evolution of gain cannot be found from the axial streak spectrum of a single shot. However long and short shots with similar coupled energies per unit length can be compared to give some indication of the growth and decay of gain. Equation (2.33) gives the axial intensity from an ASE "laser" operating below saturation as $I = S (e^{\alpha\ell} - 1)$ where $S = \epsilon/\alpha$ is the source function, ϵ the spontaneous emission coefficient and α the small signal gain coefficient. The ratio of two shots of lengths ℓ_1 and ℓ_2 is thus :-

$$\frac{\ell_1}{\ell_2} = \frac{S_1 (e^{\alpha\ell_1} - 1)}{S_2 (e^{\alpha\ell_2} - 1)} \quad (4.1)$$

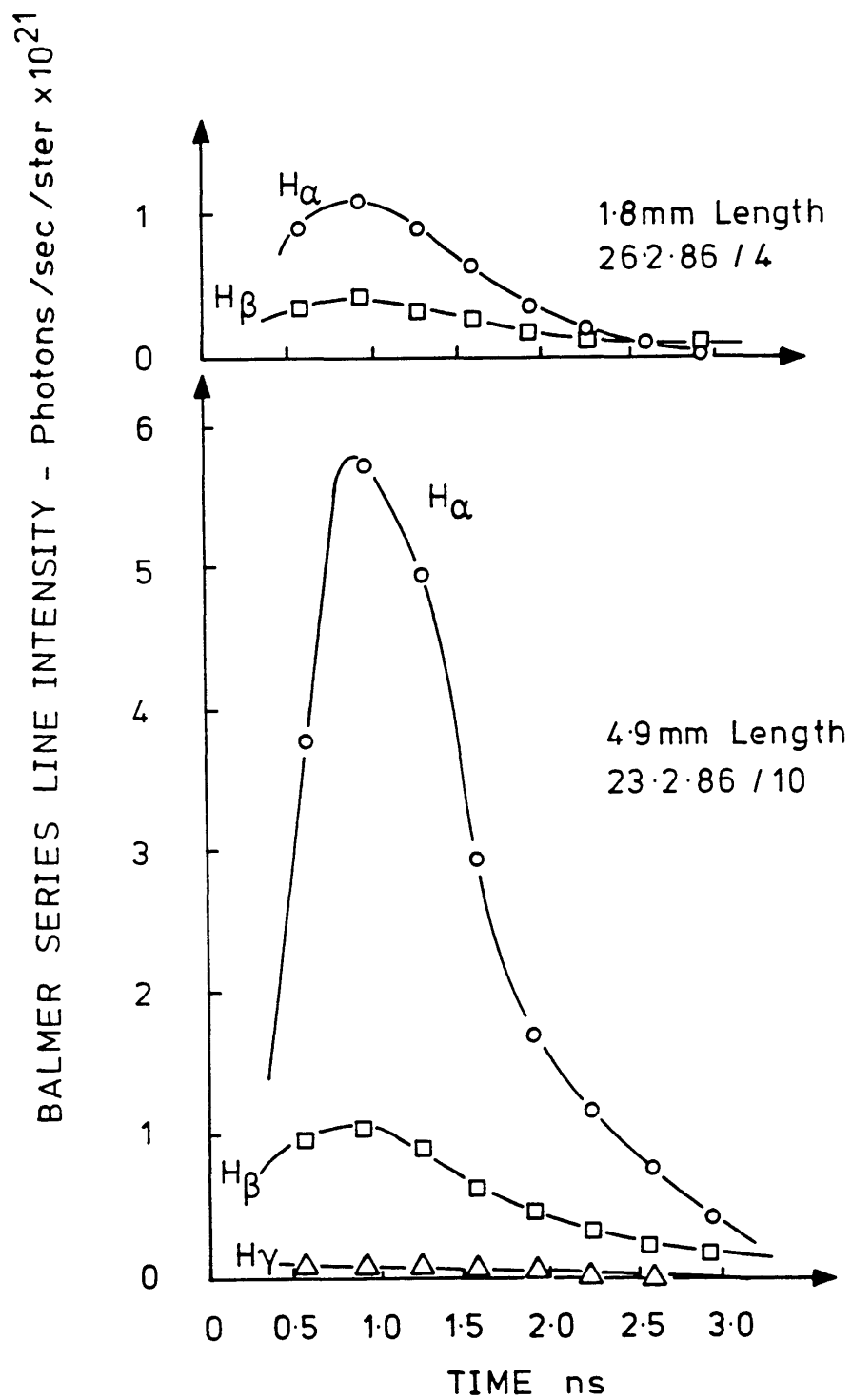


Figure 4.12 A comparison of the axial time evolution of H_{α} and H_{β} for 1.8 mm and 4.9 mm carbon fibre shots.

Clearly the source function will be a function of both time and plasma density. We do not need to know the absolute values of S_1 and S_2 , simply the ratio of the two source functions. This was estimated by Kiehn [77] for a pair of 1.9 and 6.8 mm shots by taking the ratio of the H_β lines at late time when the plasma becomes optically thin, and factoring out the lengths of the two plasmas. As the gain for H_β is negligible the ratio of the two lines gives the ratio of spontaneous emissions $\epsilon_1 \ell_1 / \epsilon_2 \ell_2$. However the spontaneous emission coefficients are simply linear functions of density and absorption oscillator strength (see equation 2.28) and so the ratio $\epsilon_1 \ell_1 / \epsilon_2 \ell_2$ should be similar for the H_α line. The ratio S_1/S_2 was estimated to be ≈ 1.47 using this method, and this allowed equation (4.1) to be solved numerically for a range of times.

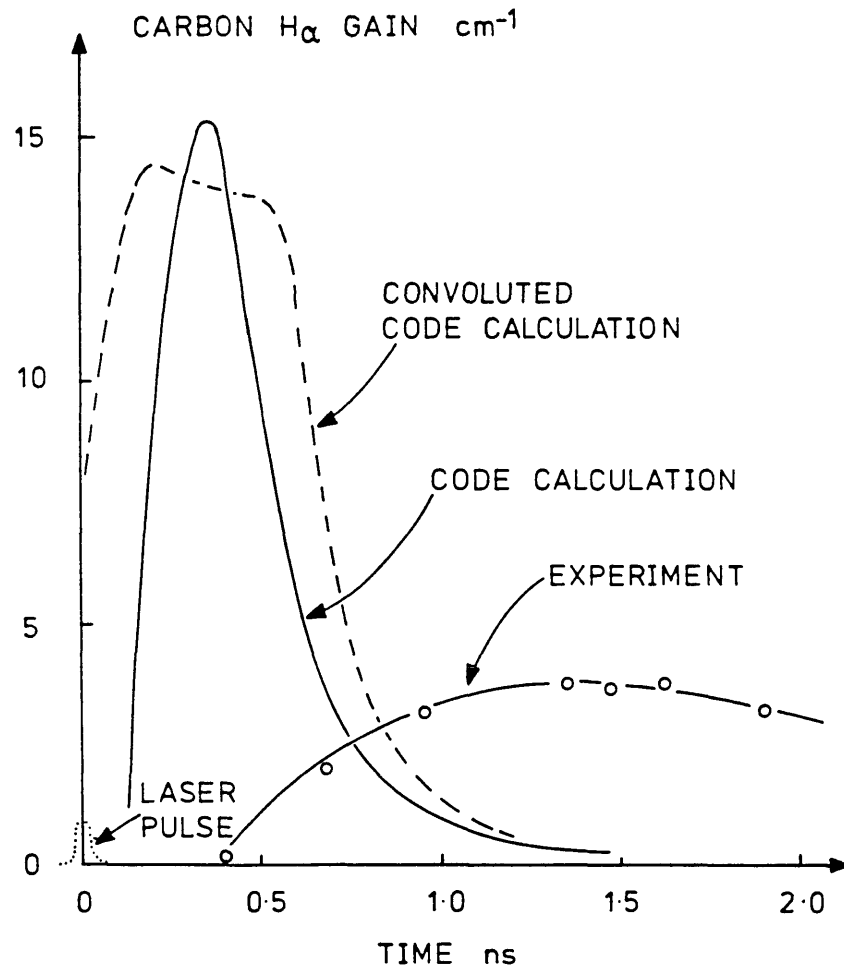


Figure 4.13 Calculated gain (with and without instrument convolution) compared with the measured time evolution of H_α gain found from axial spectra on long and short fibre shots.

Figure 4.13 compares the calculated time evolution of the gain (with the ≈ 400 ps temporal resolution of the spectrometer folded in) to the experimental gain found from comparing two axial spectra. The experimental time evolution of gain is clearly not in good agreement with computer simulations. The experimental gain peaks some 1 ns after the predicted time for maximum gain, is somewhat lower, and has a much broader temporal profile. This is believed to be due to a number of effects. Simply estimating the ratio S_1/S_2 from line intensities at one time does not take into account the time evolution of the source functions. At late times ($t > 1$ ns) the expansion of the short plasma will result in the axial and radial dimensions becoming comparable, and the plasma will continue to expand as a sphere rather than as a cylinder. The spherical expansion of the short shot enhances its cooling and gives a systematic increase in the gain found at late times.

Schlieren images of fibre targets irradiated in line focus geometry taken by Bassett and co-workers [78] have also shown fine structures in the plasma at early times. Figure 4.14 shows a time framed Schlieren image of a 3 mm long, 8 μm fibre target 500 ps after irradiation with a 20 ps, 170 J pulse of green light. These structures may result in density gradients which refract the XUV radiation we wish to amplify out of the plasma at early times. This would suppress the gain until the expansion of the plasma has become more uniform and would give lower gain as the density falls off rapidly with time.



Figure 4.14 Time framed Schlieren image of an 8 μm fibre target 500 ps after irradiation with a 20 ps, 170 J laser pulse.

4.6 SIMULTANEOUS AXIAL AND TRANSVERSE SPECTROSCOPY

In order to resolve some of the apparent disagreements between experiment and computer simulations, additional physics was incorporated in the codes and a second experiment was carried out. The experimental arrangement and method were similar to those described in 4.2. The TAE line focus was reconfigured for this experiment, such that each lens / mirror pair produced up to a 14 mm focus. Six beams were superimposed on target to improve the uniformity of irradiance.

A second time resolved XUV flat field spectrometer constructed by the author was used in addition to the axial instrument, and recorded the transverse emission from the full length of the plasma. The transverse spectrometer viewed the plasma at 35° to the axis and 7° above the horizontal. Both instruments used narrower ($\approx 200 \mu\text{m}$) streak camera timing slits to increase the temporal resolution to ≈ 80 ps, allowing the time evolution of gain to be examined in more detail. Time resolution in the transverse instrument was provided by a low magnification (0.8 x) streak camera recording over a spectral range of $175 \text{ \AA}$ at a streak rate of 105 ps / mm.

The two spectrometers were protected from plasma blow off by $1000 \text{ \AA}$ formvar filters. In addition a fine copper electron microscope mesh was used as a "neutral density" XUV attenuator on the transverse instrument. Without this mesh it was found that the high Ly_α flux this instrument received resulted in streak camera non-linearities and breakdown due to high space charge densities. The use of simultaneous axial and transverse spectroscopy allows the gain to be found on a single shot, and the time evolution of gain for one shot to be investigated without making assumption about the source function of the plasma.

To compare axial and transverse line ratios the two spectrometers must be cross calibrated. The full instrumental calibration carried out in section 3.6.5 is not sufficiently accurate for this type of investigation. However only a relative calibration between the two instruments is required and this can be done by recording the emission from a slab target irradiated in point geometry. A CH slab target was set at the usual target position with the perpendicular at 17.5° to both spectrometers, and three beams of the line focus masked down to irradiate a 1 mm section.

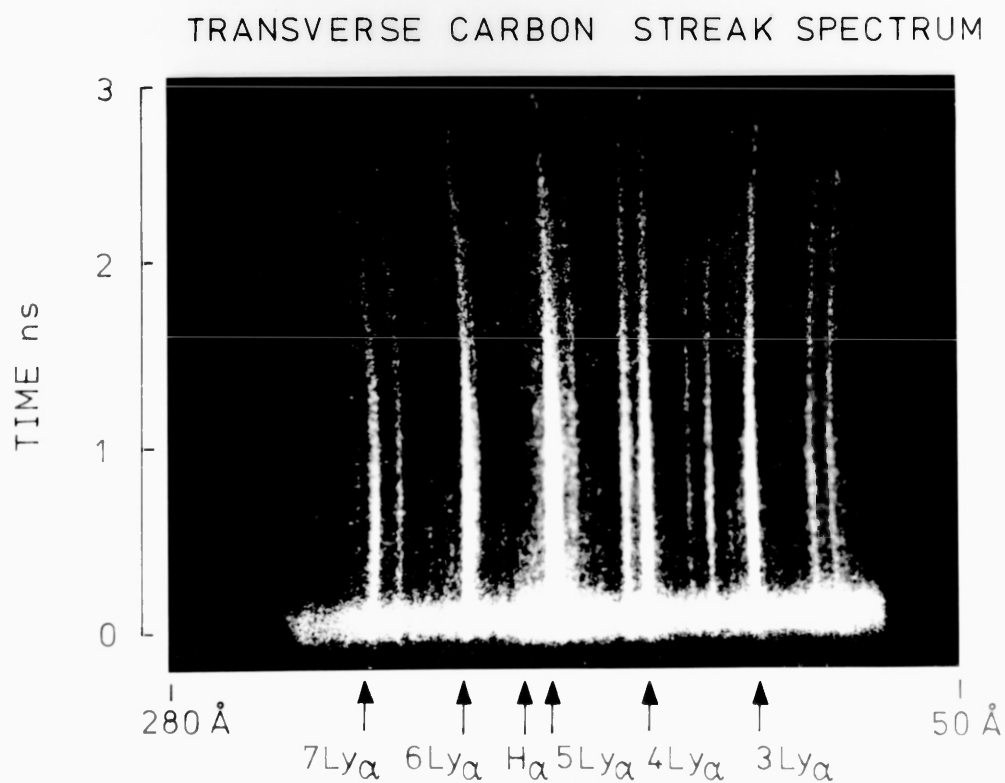
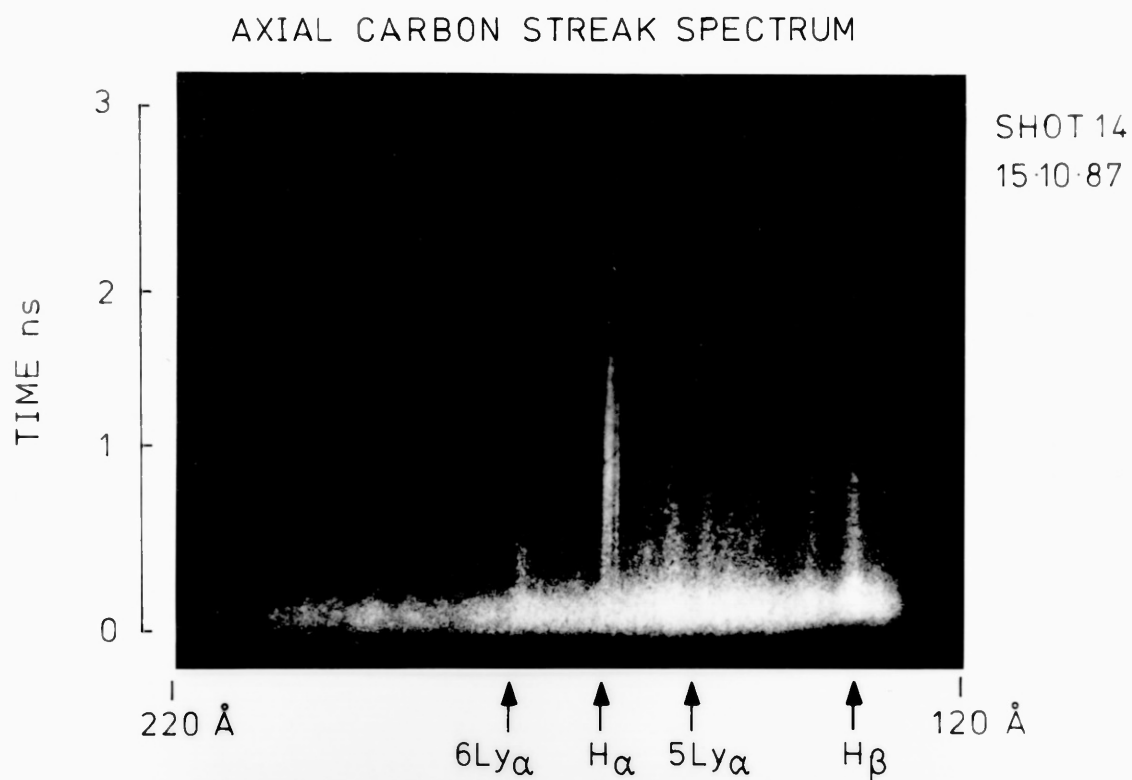


Figure 4.15 A comparison of axial and transverse streak spectra from a 12 mm carbon fibre shot.

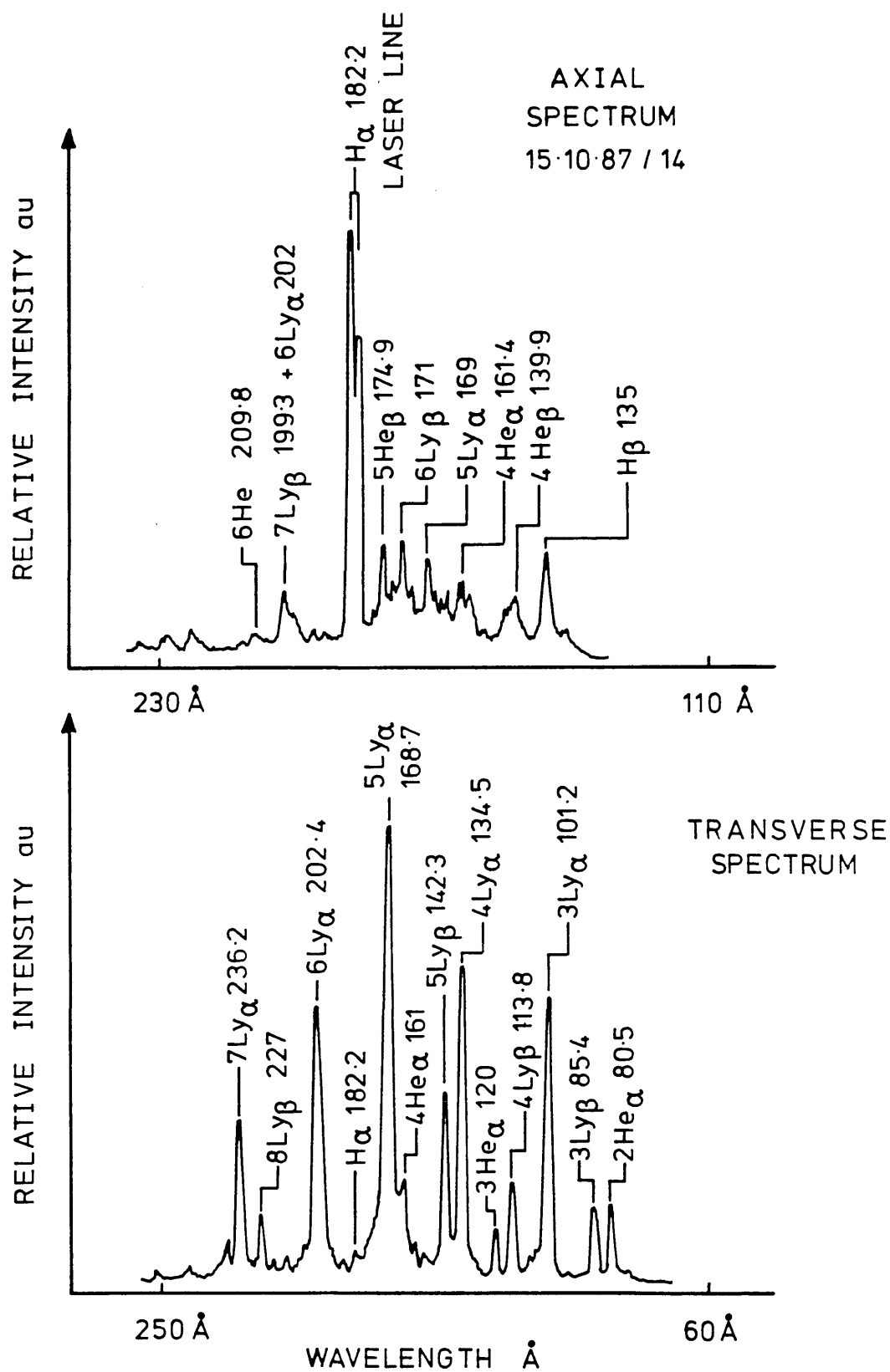


Figure 4.16 Densitometer traces (with the film curve deconvoluted) of the axial and transverse carbon fibre streak spectra presented in figure 4.15.

The predominance of resonance lines in the transverse spectrum serves to illustrate how the lower laser level in the carbon fibre scheme is efficiently depopulated. The lower state to ground transition radiation Ly_{α} must escape from the plasma, as reabsorption will clog up the population inversion. The Ly_{α} transition is optically thick in the axial direction, hence the dominance of the Balmer series in axial spectra. However in the transverse direction Ly_{α} is Doppler decoupled by the radial velocity gradient of the plasma as outlined in 2.6. This along with the small transverse dimension of the plasma prevents reabsorption of Ly_{α} in the transverse direction, leading to a population inversion by fast radiative depopulation of the $n = 2$ level.

The time evolution of gain can now be found for a single shot. The ratio of axial to transverse emission is simply :—

$$\frac{I_{ax}}{I_{trans}} = \frac{S (e^{\alpha\ell} - 1)}{\epsilon\ell} \quad (4.2)$$

$$\frac{I_{ax}}{I_{trans}} = \frac{(e^{\alpha\ell} - 1)}{\alpha\ell} \quad (4.3)$$

where ℓ is the plasma length, α the gain coefficient and ϵ the spontaneous emission coefficient. Equation (4.3) can be solved numerically for a range of times without making any assumptions about the time evolution of the source function, since the spontaneous emission coefficient ϵ factors out of equation (4.2). Figure 4.17 compares the temporal history of gain for a 2.8 mm shot with an improved simulation. The suppression of gain before 200 ps by the high temperature and density of the initial laser produced plasma is well observed experimentally, and a peak gain of 6.1 cm^{-1} is seen at some 450 ps after the heating pulse.

The gain can be seen to decay much more rapidly than was found from the axial analysis, and it is interesting to note that it occurs some 400 – 500 ps before the peak H_{α} emission. A similar effect was also seen in hydrogenic fluorine ASE "lasers" (section 5.4.4) though on a more compressed timescale, as would be expected from the Z^{-1} scaling of time in (2.60). Unfortunately the axial / transverse data set was rather limited, and a full statistical analysis of the gain from axial to transverse ratios was impossible. However this technique appears to resolve some of the disagreement between simulations and axial measurements alone.

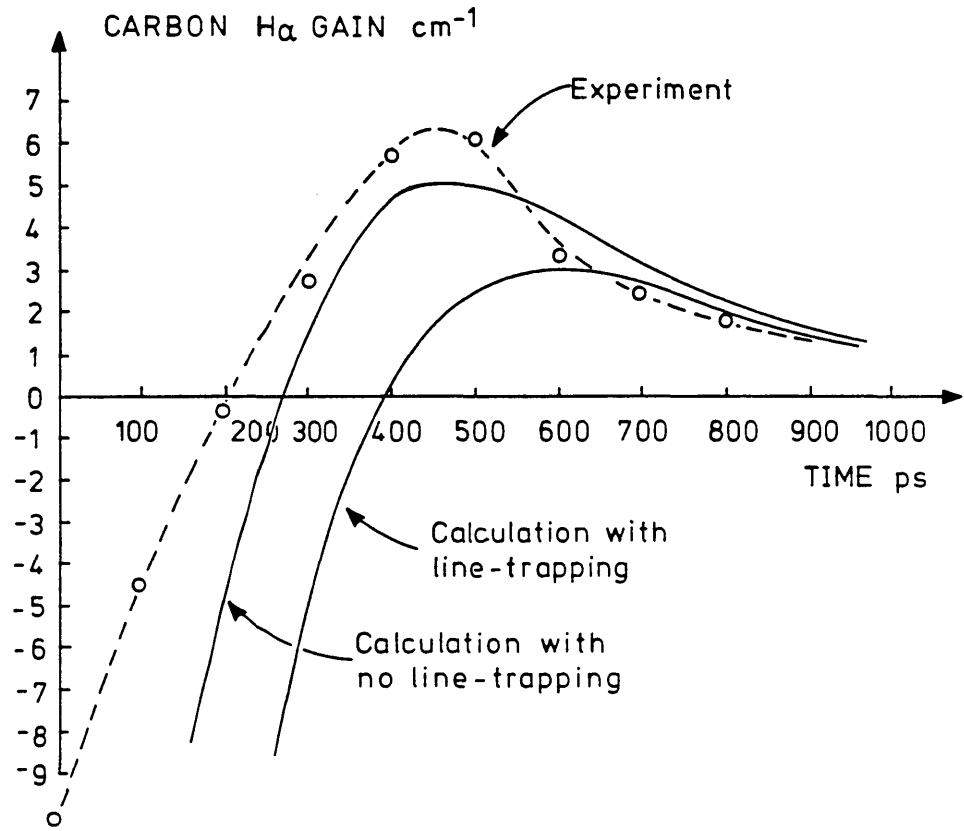


Figure 4.17 The time evolution of gain found from axial and transverse streak spectra, compared with improved code simulations (with and without line trapping).

The result of a numerical simulation of the axial transverse experiment by S. Rose [79], for a 2.77 mm long carbon fibre with 2.5 J cm^{-1} absorbed energy is shown in figure 4.17. This simulation incorporates several improvements over previous work on carbon fibres reported in section 4.5 of this chapter and recently published in reference 8. As well as correcting coding errors in the formulation of the electron-ion exchange rate and thermal conductivity in the hydrodynamic description, three improvements were made to the atomic physics content of the code.

Firstly, the modeling of the recombination and de-excitation cascade was improved by setting the upper level used in the simulation ($n=6$) to its equilibrium value. Inclusion of this term increased the calculated gain.

Secondly, the quasi-static electron Stark broadening of the H_{α} laser line was included in an approximate manner using calculations due to Oza, Greene and Kelleher [80]. Stark broadening was found to decrease the calculated gain by a factor of some 0.68.

Lastly, the fine structure splitting of H_α was taken into account. The components of H_α are sufficiently separated in energy that to a good approximation, gain on one component does not interact with gain on another. The simplest approximation of only considering gain on the $2p_{3/2}-3d_{5/2}$ component was taken. For a population distributed according to degeneracy among the fine structure levels, this component shows twice the gain of the next largest component $2p_{1/2}-3d_{3/2}$. This effect also reduces the calculated gain.

Results from two calculations are shown in figure 4.17, with and without line trapping. It appears that leaving out the line trapping gives better agreement with experiment, both in terms of the time gain turns on, and also of the magnitude of peak gain. This suggests that the local escape factor method including the Doppler decoupling used in the calculation may overestimate the Ly_α trapping. More evidence to this effect will emerge from the work on lithium fluoride coated fibres reported in the next chapter.

4.8 SPATIAL BREAK UP OF THE GAIN REGION

It is interesting to note that the source broadening of lines such as H_β can be seen in the axial streak spectrum presented in figure 4.14. This is not unreasonable as the $0.5 \text{ \AA}$ resolution of the instrument in the slitless mode corresponds to an $80 \text{ }\mu\text{m}$ source size, while the plasma is expected to expand to some $500 \text{ }\mu\text{m}$ radially within the first nanosecond. The break up of the H_α laser line into discrete elements as the annular gain region of the plasma moves away from the fibre is also visible, and this splitting was seen on a number of shots. H_α was the only line to show this effect, which was found to remain constant when the streak camera was moved slightly, showing that the split was not due to a damage spot in the photocathode.

Placing a slit on the front of the spectrometer in the same plane as the grating would enhance this observation, giving 1D space resolution in the plane of dispersion. Streaks would then record 1D spectral and 1D temporal information, with 1D spatial information extractable from the line widths.

This technique was not examined in detail on this experiment as the spectrometer was not in the correct orientation to well observe the horizontal blow out of the two plasma lobes from the fibre. However a $100 \text{ }\mu\text{m}$ slit was used on a single shot, and the linewidth of H_β used to follow the move out of the plasma boundary.

A mean radial velocity of $\approx 2.1 \times 10^7 \text{ cm s}^{-1}$ was found using this method, and the splitting of H_{α} was also used to estimate the radial position of the peak gain region as a function of time. A value of $r \approx 100 - 200 \text{ }\mu\text{m}$ was measured at $t = 500 \text{ ps}$, the time of maximum gain. Both these values are broadly in line with code predictions, but due to the limited nature of the data, no detailed comparisons were made.

4.9 CONCLUSION

The hydrogenic carbon recombination scheme was investigated using novel time resolved diagnostics at plasma lengths significantly greater than those achieved by previous workers. Gains of at least $3.35 \pm 0.33 \text{ cm}^{-1}$ at $182 \text{ }\text{\AA}$ were conclusively demonstrated in adiabatically cooled plasmas up to 9 mm in length produced by laser irradiation of $7 \text{ }\mu\text{m}$ diameter carbon fibres. It was shown that plasma lengths greater than $\approx 6 \text{ mm}$ must be investigated before gain can clearly be demonstrated by exponential growth of a line in this type of XUV laser scheme. This places a limit on the short wavelength scalability of the hydrogenic recombination scheme with our existing pump laser.

Disagreements between experimental time evolution of gain and computer simulations were partially resolved in a second more completely diagnosed experiment. Gain histories found from simultaneous axial / transverse time resolved spectroscopy gave better fits to improved simulations. This suggested that the analysis of axial data alone to find the time evolution of gain does not treat the time evolution of the source function S correctly. Using the axial to transverse technique it was also found that peak gain does not occur at the time of peak axial emission at $\approx 1 \text{ ns}$ as was assumed in the original analysis, but significantly earlier at $\approx 500 - 600 \text{ ps}$ after the heating pulse. In addition, some problems with the modelling of Ly_{α} trapping have been brought to light. A paper describing the better agreement between theory and observations of gain in carbon fibres using these improvements to both the code and experimental techniques is currently in preparation.

1D space resolution also was achieved in some time resolved axial spectra by using an imaging slit. The split of the laser line H_{α} into distinct gain regions was observed, and this allowed the time history of the radial gain profile to be found and compared to simulations.

In summary these experiments were a convincing proof of method for the H like recombination laser scheme. While $182 \text{ \AA}$ is not a particularly useful wavelength in itself, these investigations have provided a starting point for the scaling of this scheme to shorter wavelengths, which will hopefully lead eventually to a water window ASE "laser". Initial disagreements between experiment and simulations have also stimulated developments in both diagnostic techniques and computer modeling of laser plasmas, and this will result in more accurate planning and interpretation of future experiments.

CHAPTER 5. THE FLUORINE IX X-RAY LASER EXPERIMENT

ABSTRACT

The hydrogenic recombination scheme was scaled from 182 to 81 Å by changing the lasing ion from carbon [8 and chapter 4] to fluorine. Lithium fluoride coated carbon fibres were irradiated in line focus geometry and moderate XUV gain at 81 Å was observed in the freely expanding, adiabatically cooled plasmas. Three gain regimes were found for different coupled energies of 5, 9 and 16 J cm⁻¹. These were characterised by absorption (-1.8 ± 0.2 cm⁻¹), high peak gain (4.4 ± 1.0 cm⁻¹) and low peak gain (2.0 ± 1.0 cm⁻¹) on the fluorine Balmer H_α transition. This was the first observation of gain below 100 Å and the results are extremely well fitted by numerical simulations. This experiment is currently the most completely diagnosed and fully modelled X-ray laser study.

5.1 INTRODUCTION

The hydrogenic scheme may be scaled to shorter wavelength simply by using a higher Z lasing ion. The Z^{-2} scaling of wavelength is attractive to experimenters, and is accompanied by a somewhat less savage scaling of pump energy than collisional schemes.

It was decided to investigate fluorine IX plasmas as the move from carbon to fluorine gives a significant decrease in wavelength from 182 to 81 Å while gain could still be achieved in long (> 5 mm) plasmas with the available power of the VULCAN laser. To get a simple analytical estimate of the required input energy for a F IX laser we can consider the scaling of electron temperature T_e and pump laser intensity I . From a collisional radiative model [57] and an expression for the thermal limit (2.37) it can be shown that $T_e \propto Z^2$ and electron number density $n_e \propto Z^7$. From a model of laser ablation at low intensity where the absorption is dominated by inverse Bremsstrahlung [62], the electron temperature is given by $T_e \propto Z^{2/9} I^{4/9}$. It follows that the required pump laser intensity scales as Z^4 .

In the carbon laser gain was achieved at coupled energies of 2.6 J cm^{-1} and so we should expect to see gain in fluorine plasmas at coupled energies of $2.6 \times (9/6)^4 = 13.2 \text{ J cm}^{-1}$. This limits the target length to < 7 mm as to achieve this power density all six beams of the TAE line focus set up must be superimposed, three either side of the target. This does however improve the uniformity of irradiance which is an important factor in achieving maximum gain.

Targets consisted of $7 \text{ }\mu\text{m}$ carbon fibre substrates coated with $0.5 \text{ }\mu\text{m}$ of lithium fluoride (LiF). LiF was used as elemental fluorine is a highly reactive gas, while LiF is convenient for target construction as it can be vapour deposited in thin uniform layers. Lithium is a low Z ($Z = 3$) element and thus gives little contribution to the plasma mass or radiative energy loss. It also has a simple XUV spectrum and no line overlaps with the fluorine lines of interest. The $0.5 \text{ }\mu\text{m}$ LiF overcoat was thick enough to avoid burning through to the substrate and no carbon lines were seen in the spectra from these targets.

5.2 EXPERIMENTAL

This experiment was similar in concept and arrangement to the carbon laser experiment [8 and chapter 4] but was more completely diagnosed than any previous X-ray laser study at RAL. It included the first use of simultaneous axial and transverse time resolved spectroscopy (Figure 5.1).

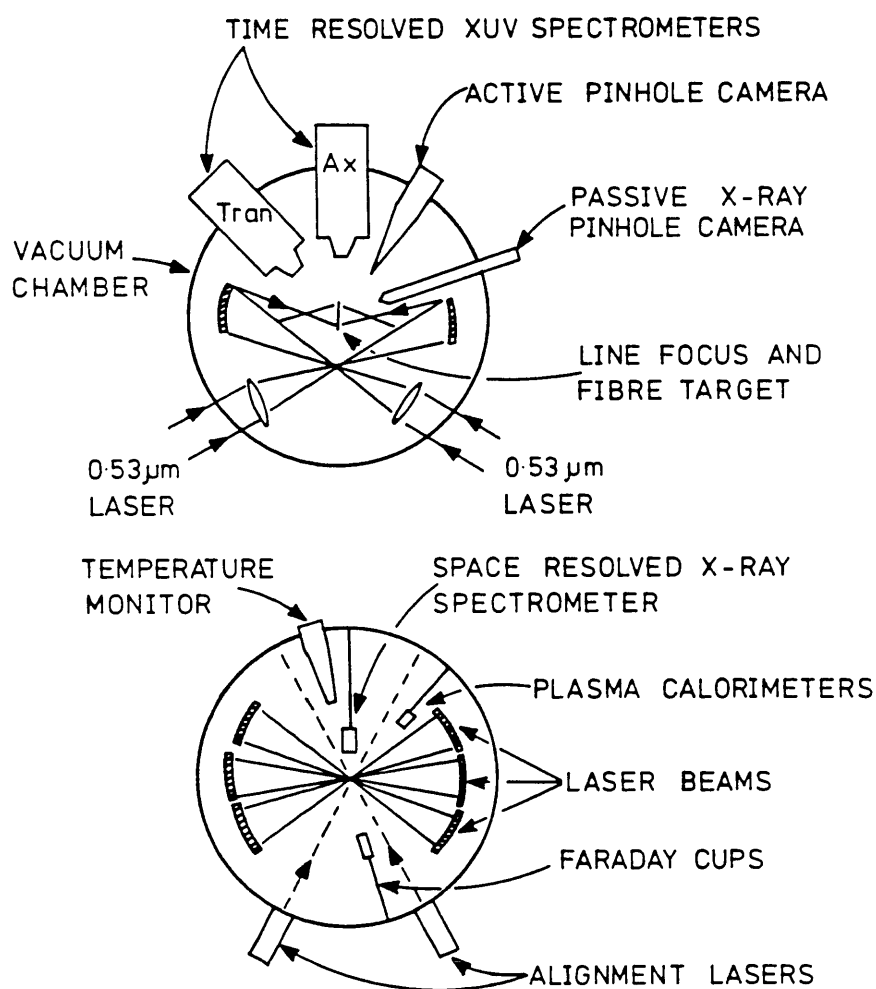


Figure 5.1 Experimental arrangement for the fluorine X-ray laser study.

$8\mu\text{m}$ diameter LiF fibre targets supported at one end were irradiated with 70 ps pulses of $0.53\mu\text{m}$ green light from the VULCAN laser at RAL. Fibre targets were aligned to $\pm 5\mu\text{m}$ in position and 10^{-3} rad in angle using a split field microscope system [68]. The six beams of the line focus were individually aligned on target using a CW argon ion laser, and for lengths less than 7 mm some of each beam masked off.

A pair of time resolved flat field 1200 lines / mm XUV spectrometers [14 and chapter 3] were used to observe line emission over a 40 – 140 Å spectral window with 0.5 Å spectral and 80 ps temporal resolution, over a 3 ns time window. One instrument was aligned along the fibre axis while the other viewed the full length of the plasma at 35° to the axis and 7° above the horizontal. Both instruments were protected from plasma blow off by 4000 Å thick formvar filters. The use of two such spectrometers allowed shot to shot variations in pump laser energy to be taken into account, and single shot measurements of gain as a function of time to be made. Plasma length was recorded with active and passive pinhole cameras filtered to accept $\approx$ 500 eV photons. Coupled energy was measured with arrays of plasma calorimeters and Faraday cups. Plasma uniformity along the fibre axis was investigated using an active temperature monitor comprising a space resolving slit, beryllium filter pack, X-ray phosphor and image intensifier. The F IX resonance spectrum at $\approx$ 15 Å was space resolved transverse to the fibre axis using a 1D penumbraally imaging X-ray crystal spectrometer to give a radial profile of the hydrogenic population density.

A comprehensive data set was collected for a range of target lengths from 1.5 to 5.8 mm and coupled energies from 4 to 20 Jcm⁻¹. The range of energies in the data set allowed the plasma energy optimisation of the F IX laser scheme to be studied. X-ray pinhole images showed the plasmas produced to be very uniform, and the temperature monitor gave an average plasma temperature of 250 $\pm$ 100 eV close to the peak of the heating pulse. While this diagnostic was time integrated, most of the emission it saw came from the plasma at early time when the density was high.

5.3 CROSS CALIBRATION OF XUV SPECTROMETERS

The primary diagnostics on this experiment were the pair of time resolved XUV spectrometers. To be able to compare axial and transverse line intensities these two instruments must be cross calibrated. The calibration described in Chapter 3 is inaccurate below 100 Å and small variations in individual photocathodes and filters must be taken into account.

The cross calibration was carried out with a slab target consisting of 5 μm of LiF coated onto a glass slide. The target was set in the usual target position with the perpendicular at 17.5° to both spectrometers. A 0.5 mm length of the target was irradiated at the power level used for fibre shots. The axial and transverse spectra were similar (Figure 5.2) and consisted primarily of He and Li like rather than H like lines. This was thought to be due to lateral heat transport across the slab surface resulting in a cooler than expected plasma. The time evolution of various lines was examined to find the peak of emission and this was used as a temporal reference between the two spectra which were recorded at slightly different sweep rates. The axial / transverse ratio for several lines was found for a range of times and an average over 500 ps was calculated. Figure 5.3 shows the axial / transverse ratio as a function of wavelength.

SOLID LiF CROSS CALIBRATION SPECTRA

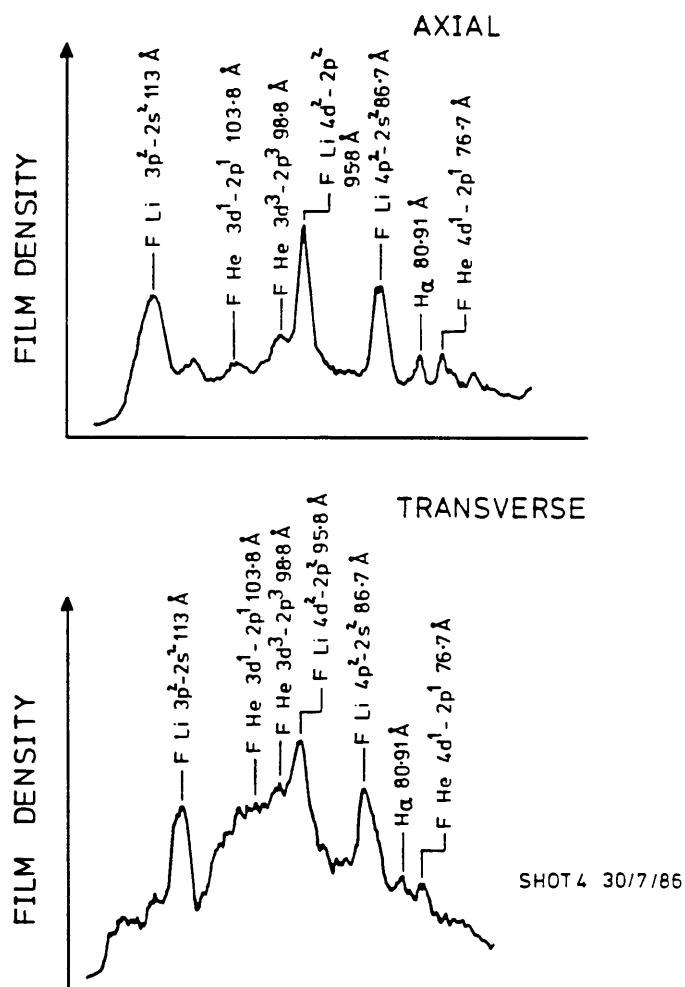


Figure 5.2 Axial and transverse cross calibration spectra from a LiF slab.

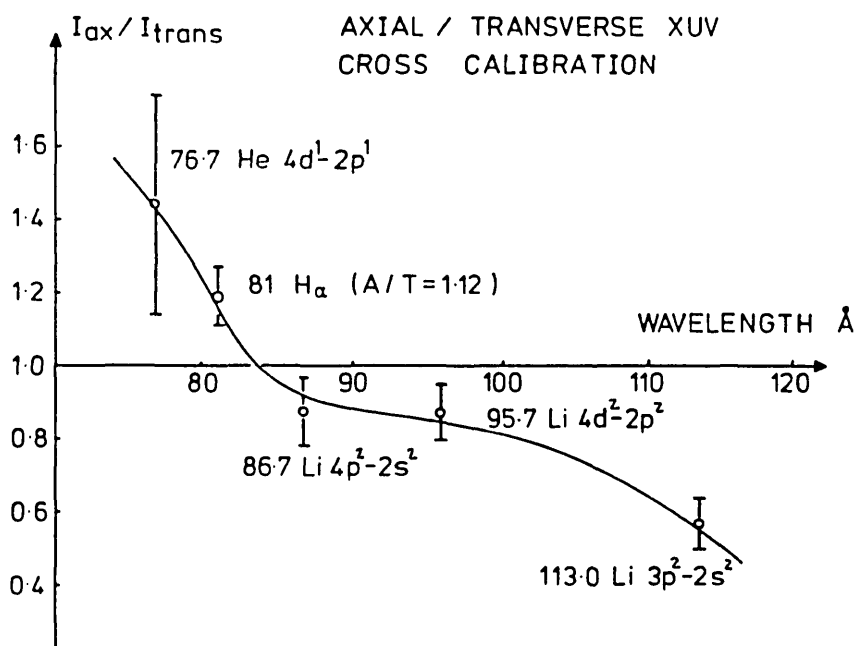


Figure 5.3 Axial / transverse spectrometer cross calibration curve.

The variation of relative sensitivity with wavelength of the two spectrometers is thought to be due to differences in the coating thicknesses of the CsI photocathodes, slight variation in filter thicknesses and differences in the efficiency of the two gratings. Error bars are due to statistical fluctuations in the axial / transverse ratios at various times and arise from the uncertainty in the level of background signal on which the spectra are superimposed.

5.4 EXPERIMENTAL RESULTS.

5.4.1 XUV SPECTRAL OBSERVATIONS

Figure 5.4 shows a time integrated spectrum from a solid LiF shot taken with the transverse spectrometer. As with time integrated carbon spectra [Chapter 4] this spectrum is again dominated by many high order H and He like resonance lines. A somewhat surprising result is that the 1200 lines / mm grating appears to still disperse at $14.98 \text{ \AA}$ (F IX Ly_{α}). Again the proposed laser line F IX Balmer H_{α} at $80.91 \text{ \AA}$ is barely resolved and is the only member of the Balmer series visible. The predominance of short wavelength higher order resonance lines in this spectrum is due to the higher sensitivity of the phosphor plate detector at shorter wavelength (see section 3.6.3).

SOLID LiF TIME INTEGRATED XUV SPECTRUM

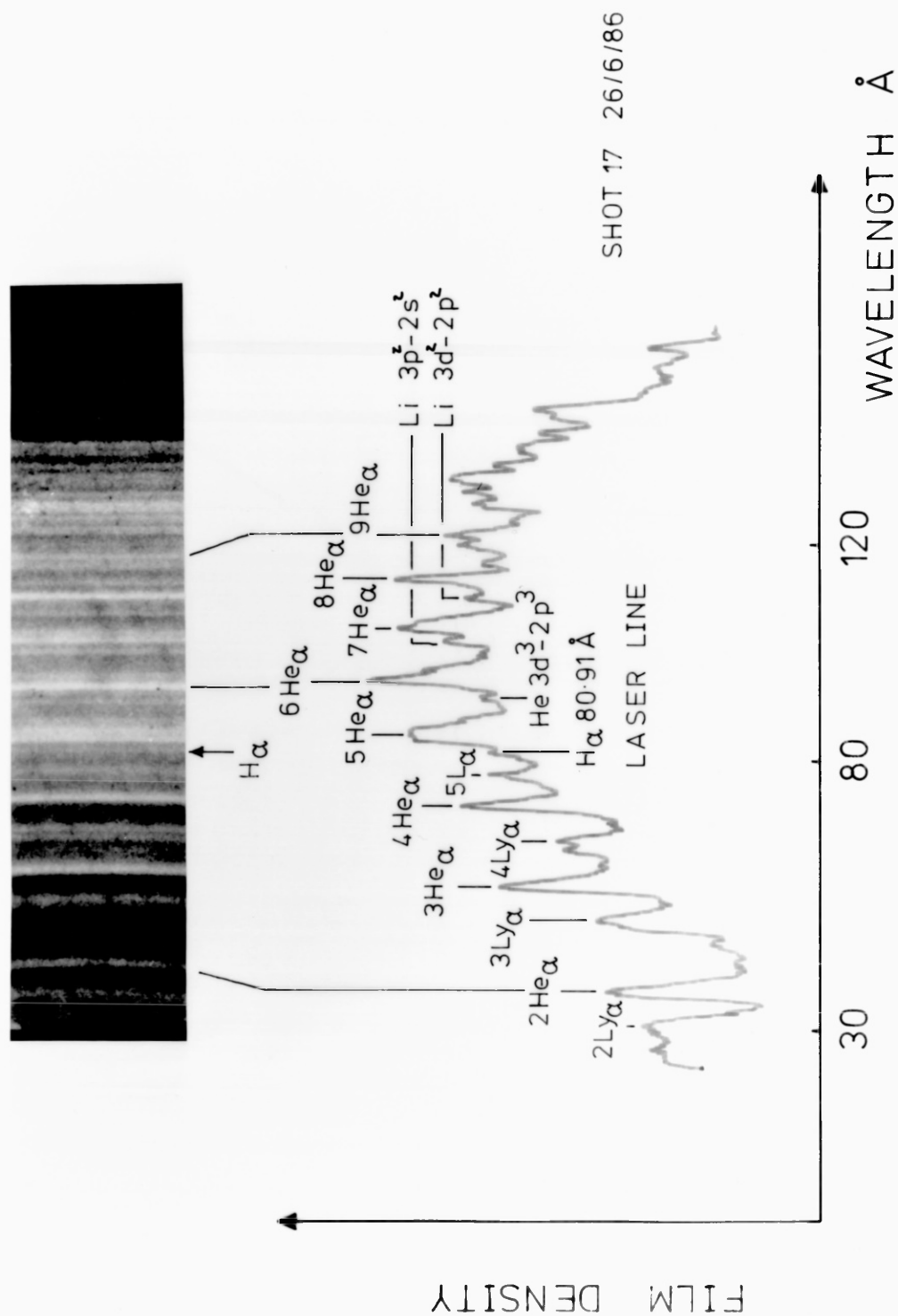


Figure 5.4 A time integrated transverse XUV spectrum from a LiF coated fibre target.

Figure 5.5 compares axial and transverse streak spectra from a single 5.5 mm long LiF fibre shot with $\approx 10 \text{ J cm}^{-1}$ coupled energy. These spectra are quite different and distinctive, clearly demonstrating anisotropy between axial and transverse emission. Comparing traces (film curve deconvolved) of these spectra taken at 250 ps after the start of the laser pulse (Figure 5.6) it can clearly be seen that the axial spectrum is dominated by the Balmer series and that H_{α} is by far the most prominent feature. The transverse spectrum however is composed of helium like and hydrogenic lines of comparable intensity. This result is expected, as optically thick lines will be suppressed in the axial spectrum while lines with gain will be enhanced. Indeed the axial to transverse ratio is our primary indicator of gain.

The transverse spectrum is also broadened compared to the axial spectrum. The plasma expands away from the fibre target as two lobes in the horizontal plane. The axial spectrometer views in the horizontal plane with the grating also in the horizontal plane. Source broadening in the axial spectrum is thus due solely to expansion in the vertical direction. The transverse spectrometer however views the plasma at 7° to the horizontal and consequently a component of the expansion in the horizontal plane also contributes to the transverse source broadening.

The transverse spectrometer also sees a radial Doppler component which the axial instrument does not. Indeed radial Doppler decoupling of Ly_{α} is essential for the operation of this laser scheme. However assuming a radial expansion velocity of $\approx 10^7 \text{ cm s}^{-1}$ this only contributes $\approx 0.06 \text{ \AA}$ to the width of an $81 \text{ \AA}$ line, and this is well below the resolving power of the instrument.

5.4.2 EVIDENCE FOR MULTIPLE PLASMA ENERGY REGIMES

Both length and plasma energy variations are included in the data set and these must be characterised before looking for evidence of gain. Figure 5.7 shows a plot of transverse H_{α} brightness (arbitrary units) at 250 ps against plasma length. Plasma length L_{90} was taken to be that from which 90 % of the 500 eV photons in the APC image were emitted. As the plasma is only a few hundred microns thick in the transverse direction and should be optically thin for the Balmer series, transverse H_{α} intensity should vary linearly with plasma length.

XUV STREAK SPECTRA FROM A LiF FIBRE TARGET

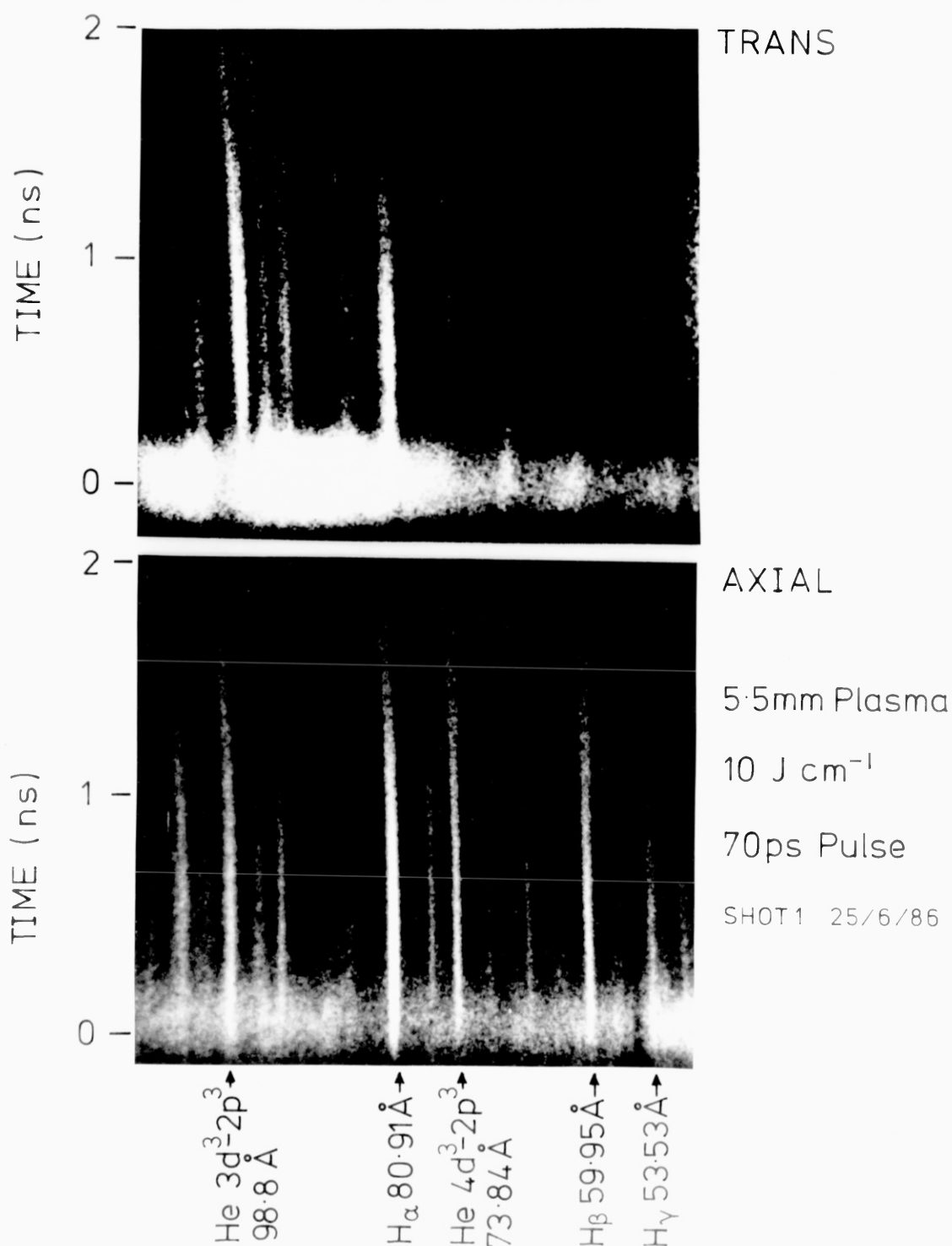
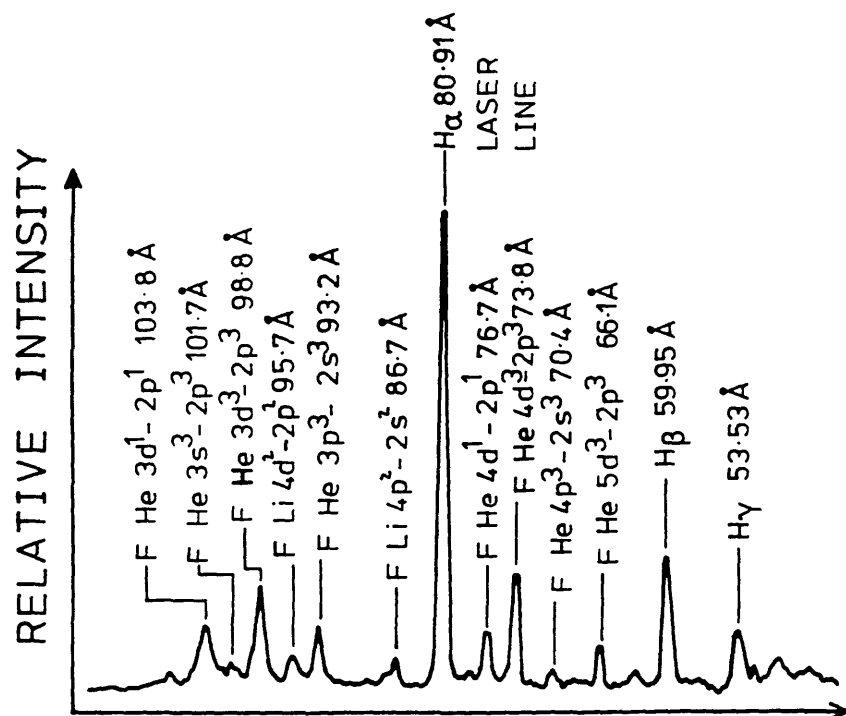


Figure 5.5 Axial and transverse XUV streak spectra from a 5.5 mm long LiF coated fibre target.

AXIAL FLUORINE SPECTRUM



TRANSVERSE SPECTRUM

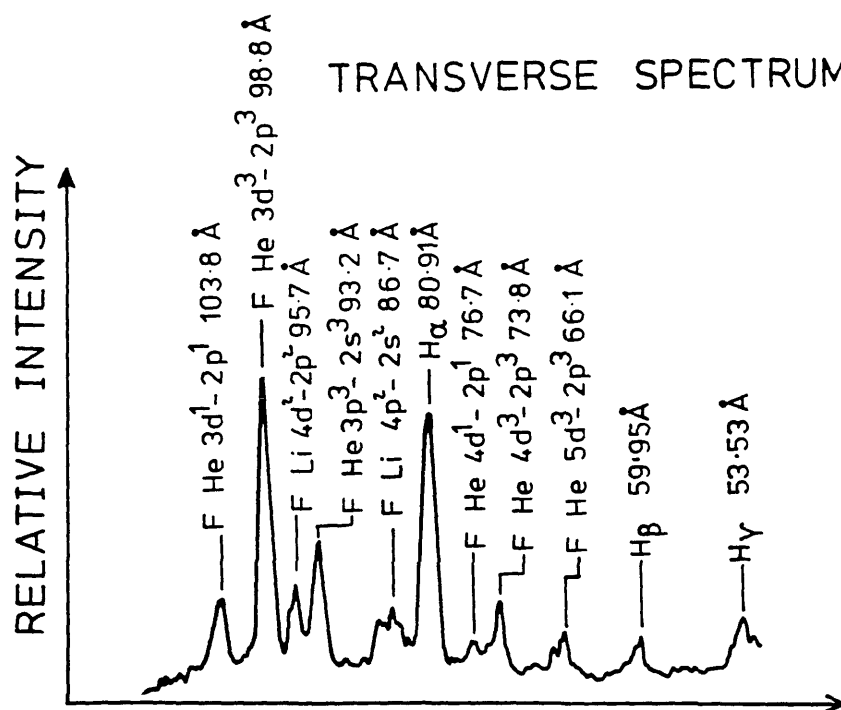


Figure 5.6 Densitometer traces of the spectra presented in figure 5.5, taken 250 ps after the start of the laser pulse.

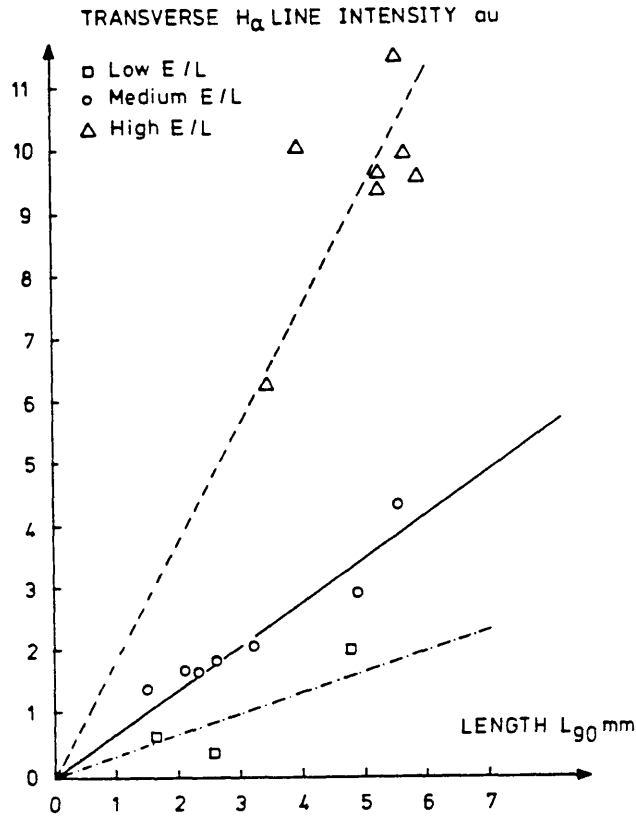


Figure 5.7 Transverse H_{α} brightness versus length, showing splitting into three distinct regimes.

What is actually seen is that there are three distinct linear regimes which are characterised by low, medium and high transverse H_{α} brightness. In the arbitrary intensity units of figure 5.7, the brightnesses / mm of these regimes are 2.00 ± 0.1 , 0.76 ± 0.09 and 0.3 ± 0.1 . The three regimes are separated from each other by more than 4σ and so the split is statistically highly significant. Similar and consistent results are found for H_{α} at a range of later times, and also for transverse H_{β} brightness over a range of times.

The split into three distinct transverse brightness regimes indicates that the data set contains contributions from several discrete values of plasma energy per unit length (E/L). Figure 5.8 is a plot of total plasma energy from Faraday cups and calibrated plasma calorimeters versus plasma length. It can clearly be seen that there are again three regimes of absorbed energy, and these are consistent on a shot to shot basis with the results from figure 5.7. We shall call these regimes "low E/L ", "medium E/L " and "high E/L " for convenience. They are characterised by absorbed energies of $16 \pm 1 \text{ J cm}^{-1}$, $9 \pm 2 \text{ J cm}^{-1}$ and $5 \pm 0.7 \text{ J cm}^{-1}$. Again the split between regimes is statistically significant at $> 2\sigma$.

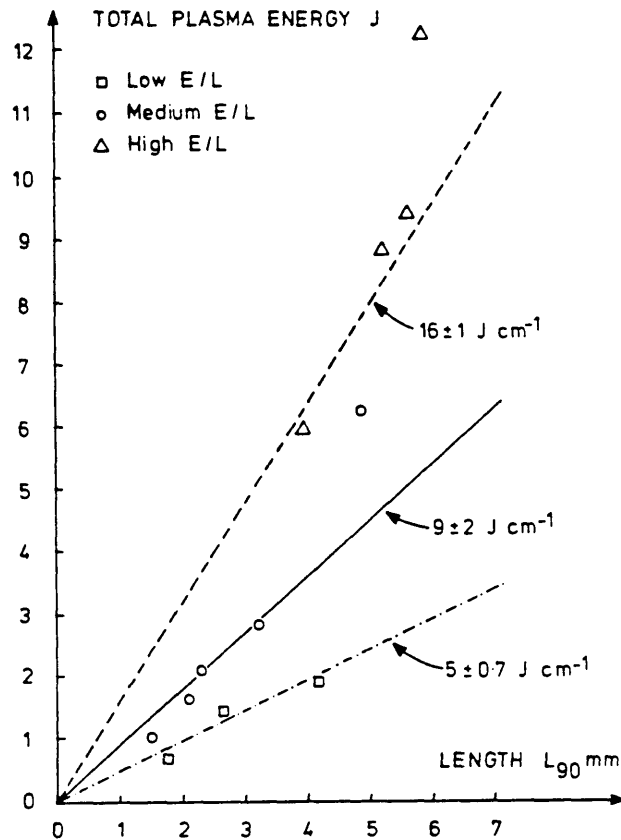


Figure 5.8 Total plasma energy versus plasma length, showing splitting into three distinct regimes.

The total pinhole camera image brightness is also a good qualitative indicator of plasma energy. As expected, a plot of total brightness versus length (Figure 5.9) shows splitting between the high and medium E/L regimes, consistent with figures 5.7 and 5.8. The splitting of the low E/L regime is however less obvious in this data set.

We should expect obvious spectral differences between the three regimes as the variation in plasma energy will result in a different ionisation balance for each regime. Figure 5.10 compares transverse XUV spectra taken at 250 ps for ≈ 3 mm shots from the low, medium and high E/L regimes. The high E/L spectrum contains H like lines (H_{α} and H_{β}) and He like lines ($3d^3D - 2p^3P$, $3d^1D - 2p^1P$ and $3p^3P - 2s^3S$) of comparable intensity. In the medium E/L spectrum He like lines predominate though H_{α} is still one of the brighter lines in the spectrum, while the low E/L spectrum is completely dominated by He like lines.

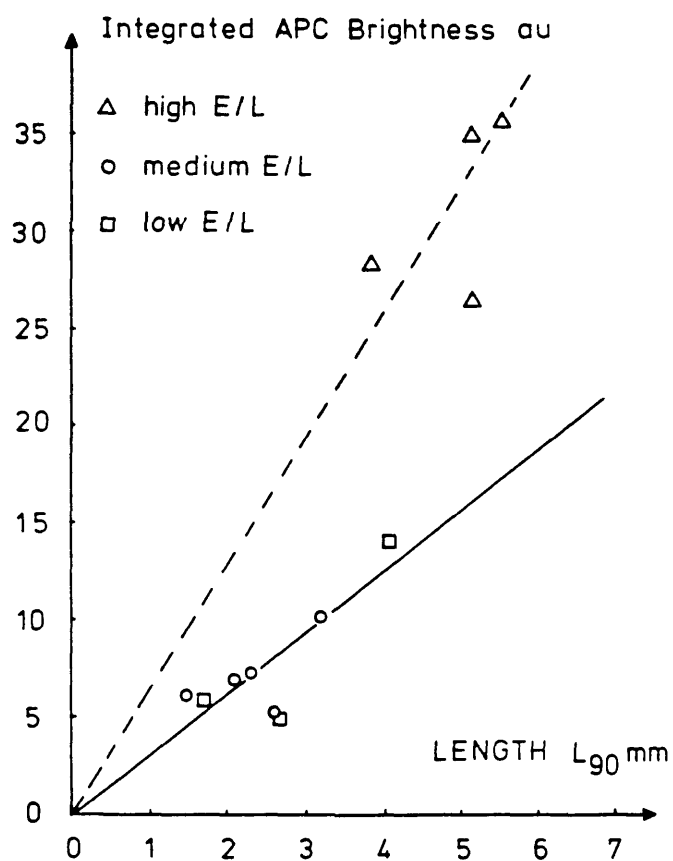


Figure 5.9 Total pinhole camera brightness versus plasma length.

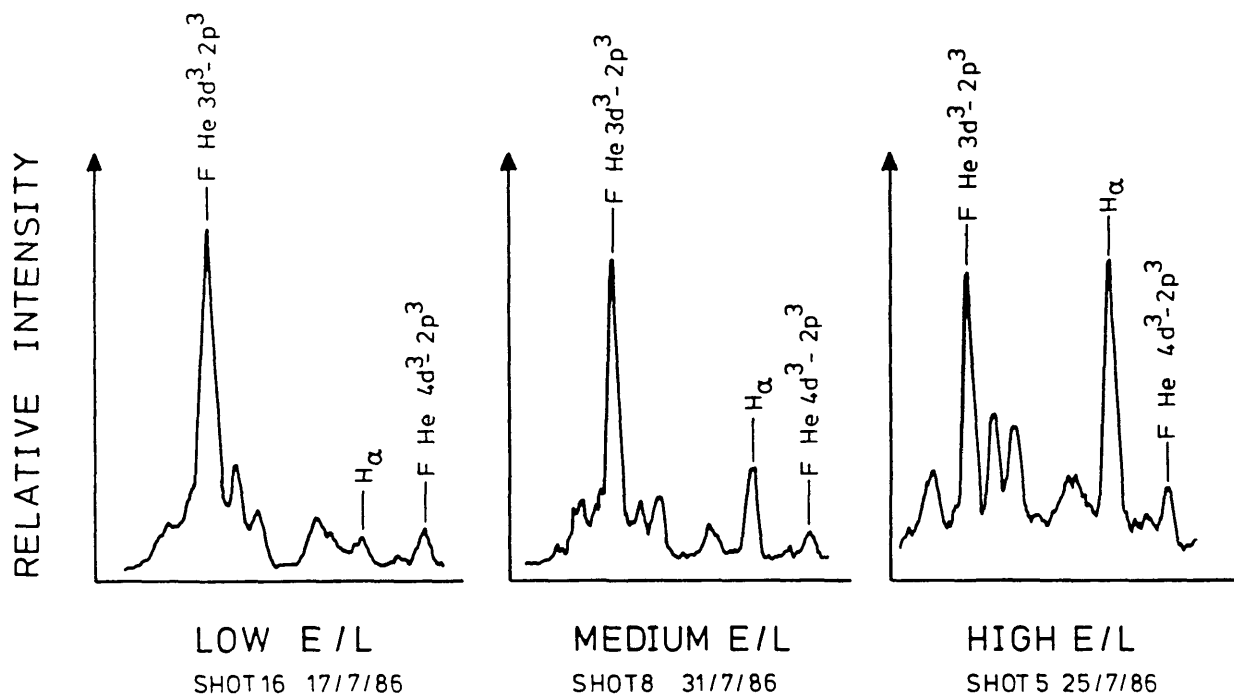


Figure 5.10 Transverse spectra from low, medium and high E/L shots of similar length.

These spectral differences are summarised in figure 5.11 which compares the ratio of hydrogen and helium like lines ($H_{\alpha} / \text{He } 4d^3D - 2p^3P$) for the three regimes plotted against plasma energy / unit length. Again a splitting into three broad regimes can be seen, and these results are consistent with the assignment of regimes from the transverse H_{α} brightness.

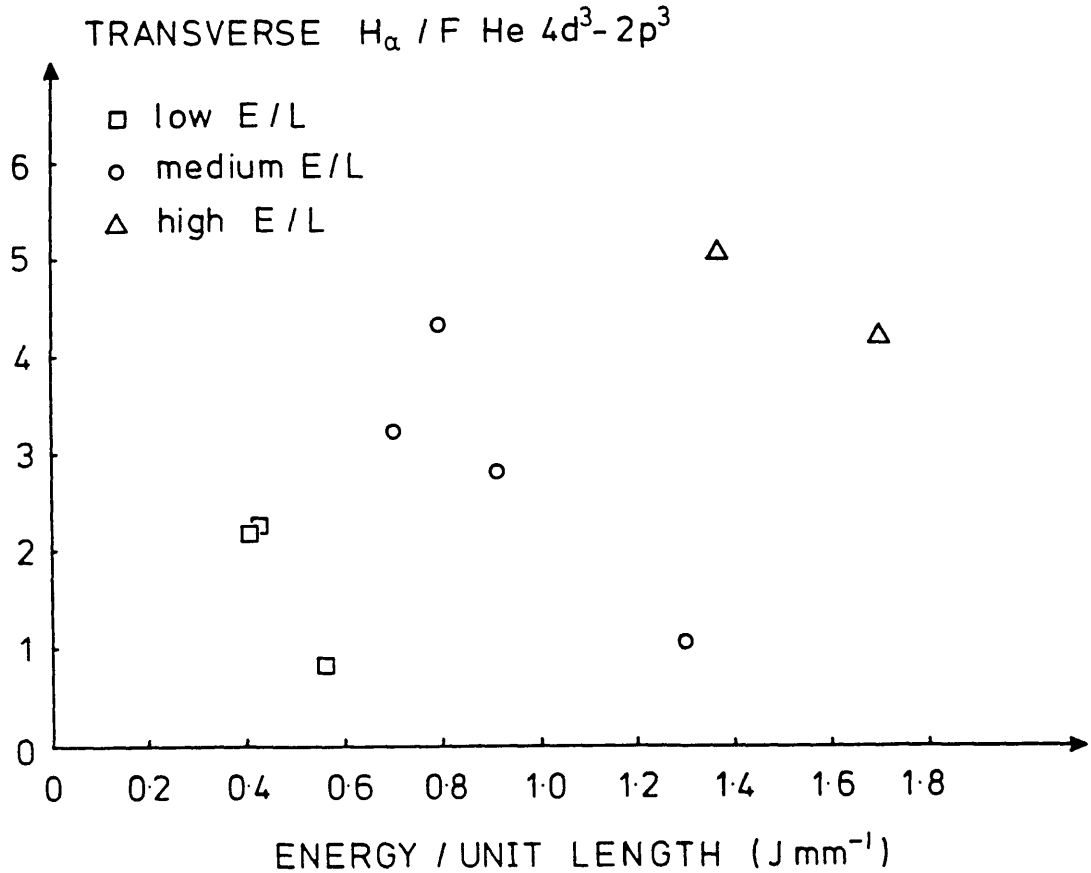


Figure 5.11 The ratio of transverse H_{α} to $\text{He } 4d^3-2p^3$ versus plasma energy per unit length showing splitting into three broad regimes.

5.4.3 OBSERVATION OF GAIN

In section 5.4.2 it was shown that three distinct regimes of absorbed plasma energy are represented in the data set, and that they are also characterised by different ion populations. The matching of plasma mass, energy and expansion rate is critical for producing gain in this type of scheme [Chapter 2]. For this reason we will treat these three regimes as separate entities when looking for evidence of gain.

Our main indicator of gain is the axial / transverse ratio variation with plasma length. For a line with gain or attenuation coefficient α , the axial / transverse line ratio is given by $I_{\text{ax}} / I_{\text{trans}} = (e^{\alpha \ell} - 1) / \alpha \ell$. This presumes that any gain does not saturate and significantly effect the radiation field, and that opacity effects are the same in the axial and transverse directions. Figure 5.12 is a plot of $H_{\alpha} \text{ ax} / H_{\alpha} \text{ trans}$ versus length for the three regimes identified in section 5.4.2. As expected the three regimes fall onto three distinct growth curves. Curve fitting with a least squares method to a function of the form $(e^{\alpha \ell} - 1) / \alpha \ell$ gives gain coefficients of $2.0 \pm 1.0 \text{ cm}^{-1}$ and $4.4 \pm 1.0 \text{ cm}^{-1}$ for the high and medium E/L regimes and an absorption coefficient of $-1.8 \pm 0.2 \text{ cm}^{-1}$ for the low E/L regime.

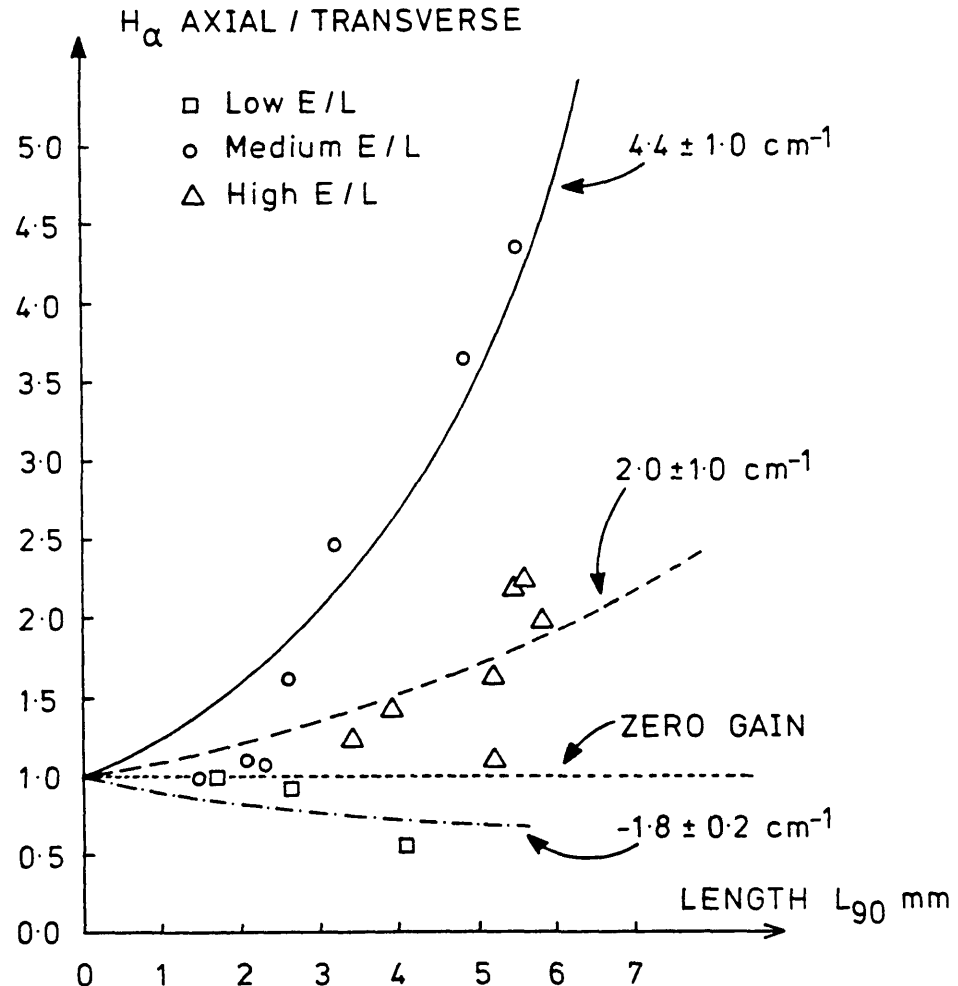


Figure 5.12 A plot of axial / transverse H_{α} versus plasma length, for the three energy regimes identified in section 5.4.2.

The lower gain observed in the high E/L shots (even though they have a higher H like ion population) suggests that either the plasma is too hot, inhibiting population of the $n = 3$ upper laser level at early times, or that the increased H like ion population increases Ly_{α} trapping more than it increases the gain. In the low E/L case the plasma is too cold, and while there is still a significant population of H like ions, the high H like ground state population and lower radial expansion velocity lead to increased Ly_{α} trapping. This prevents the inversion of the $n = 3 - 2$ levels and the plasma becomes optically thick for H_{α} .

As in chapter 4, gain may also be found from the axial intensity variation with plasma length. The axial intensity can be fitted to a growth curve of the form $I_{ax} = S.(e^{\alpha\ell} - 1)$. Using a log / linear least squares fit to the axial data gives similar and consistent values of gain for the medium and high E/L regimes, while the low E/L regime gives a good fit to a linear growth curve. Figure 5.13 shows the fit to the axial H_{α} intensity alone. Note that the split between the medium and high regimes is much less obvious than for the axial to transverse ratio method.

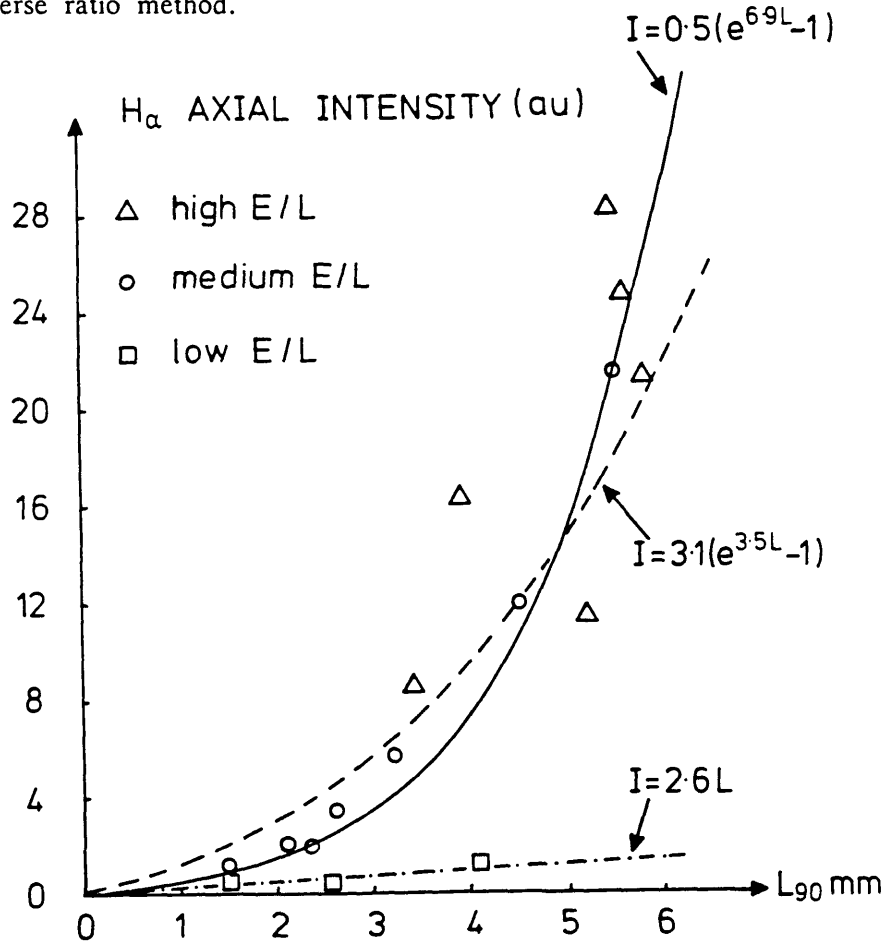


Figure 5.13 Growth curves for axial H_{α} intensity with plasma length.

This is due to the appearance of the source function $S = \epsilon/\alpha$ (ϵ the spontaneous emission brightness, α the gain coefficient) in the axial growth equation, and this can cause problems when trying to compare shots with different coupled energies. While the high E/L shots have lower gain their level of spontaneous emission is higher as they are hotter, and consequently they have a larger source function. The high and medium regimes are however distinguished by the much sharper rise in intensity of the medium E/L regime at longer lengths.

A similar analysis was attempted for the Balmer H_β and H_γ lines, which are inverted, but for which no significant gain was expected as the oscillator strengths for these transitions are too low. Figure 5.14 shows a plot of axial H_γ brightness at 250 ps against length, which to a good approximation is linear for the three E/L regimes, indicating no gain. A full axial / transverse analysis of these lines was however impossible due to the lack of a cross calibration in this wavelength range, and the large signal to noise ratio.

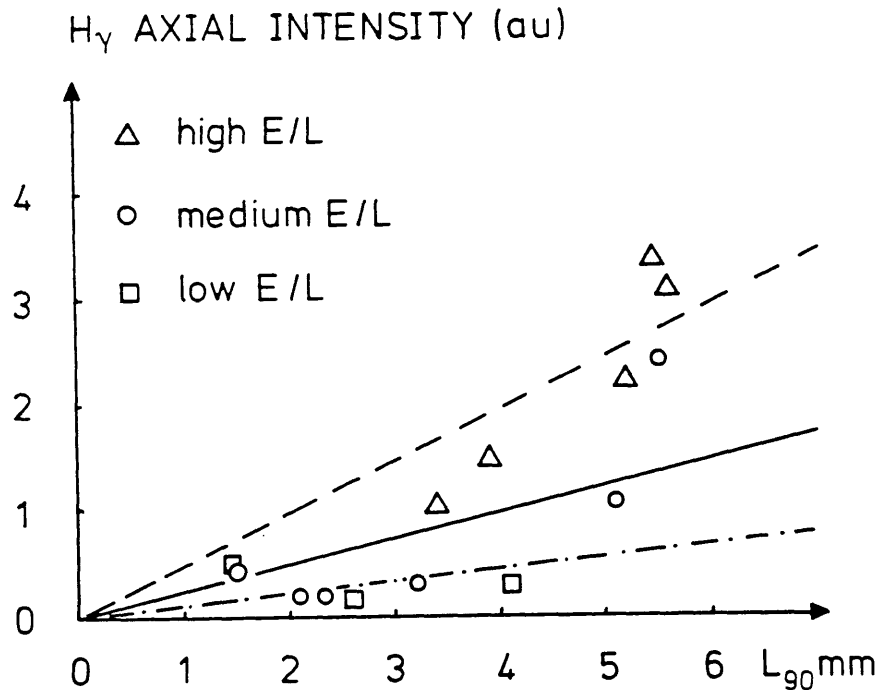


Figure 5.14 Linear growth curves for axial H_γ intensity.

5.4.4 TIME EVOLUTION OF GAIN

One of the major advantages of simultaneous axial and transverse spectroscopy is that the time evolution of gain can be found from a single shot. The comparison of axial brightnesses from long and short length shots as carried out in chapter 4 does not treat the time evolution of the source function correctly, as ϵ the spontaneous emission signal is a strong function of time, and this can lead to significant errors at late times.

By solving $I_{ax} / I_{trans} = (e^{\alpha \ell} - 1) / \alpha \ell$ numerically for one shot over a range of times, a more accurate picture of the time evolution of gain can be built up. Figure 5.15 compares the time evolution of axial and transverse emission with gain for a 5.5 mm long medium E/L shot. It is interesting to note that peak gain is not at the time of peak axial emission as was assumed in earlier work on carbon [8 and chapter 4]. As with carbon fibres, peak gain is significantly earlier than peak axial emission, at 250 rather than 400 ps after the laser pulse and results from the later turn on of the transverse as compared to axial emission.

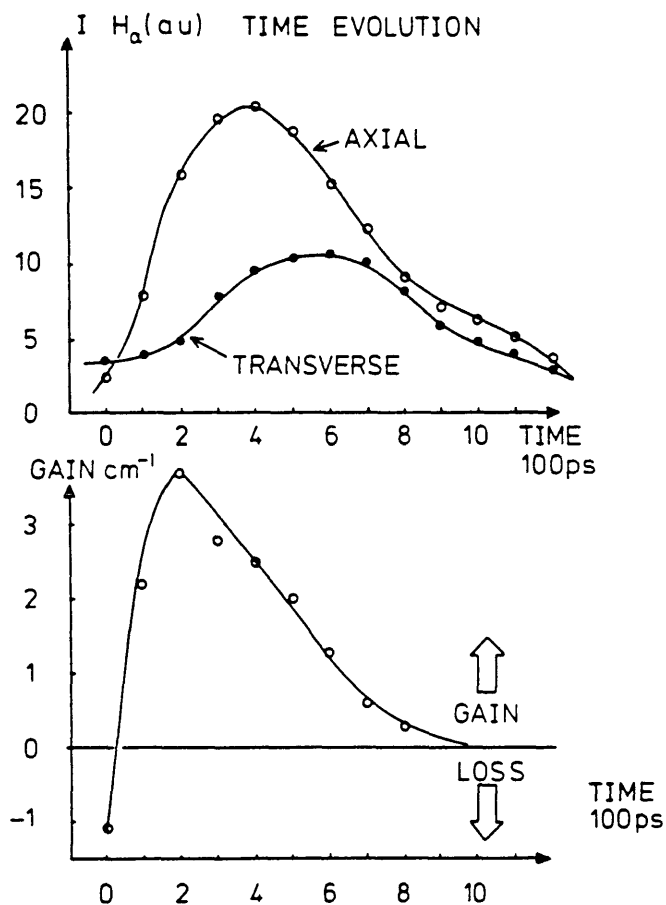


Figure 5.15 A comparison of the axial and transverse emission of H_{α} and the measured gain as a function of time.

5.4.5 OBSERVATIONS SUPPORTING GAIN

The unambiguous axial / transverse observations of gain reported in 5.4.3 and 5.4.4 are supported by several additional results which in themselves do not conclusively demonstrate gain, but are consistent with gain on H_{α} . Clear evidence of population inversion on the $n = 3-2$ levels is seen in the data from the 1D space resolving spectrometer. Figure 5.16 shows the $n = 2, 3$ and 4 hydrogenic populations as a function of radial distance from the target for a shot in the medium E/L high gain regime. The population inversion develops in space as the plasma expands away from the target, cools and recombines. It is interesting to note that the $n = 4-2$ levels are also inverted. No significant gain is however expected for this transition as the oscillator strength is too low. Few data shots were recorded with this instrument and unfortunately it is impossible to compare the population profiles for the three E/L regimes.

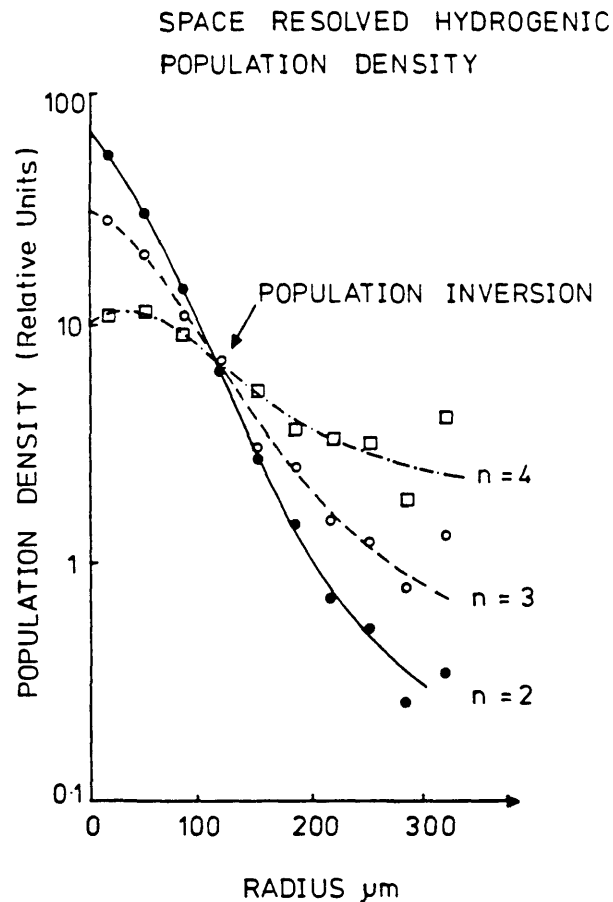


Figure 5.16 Radial profile of the $n = 2, 3$, and 4 levels showing the development of the population inversion in space.

Support for gain on H_{α} also comes from the comparison of axial to transverse ratios for lines other than the laser line. Figure 5.17 shows the axial to transverse ratios for a number of lines at 250 ps taken from shots in the three E/L regimes. Clearly H_{α} at 81 Å is the only line with an axial to transverse ratio significantly greater than 1, and it is only greater than 1 for shots in the medium and high E/L regimes. All other lines have ratios less than or equal to 1 and are optically thin with no gain (He $4d^1-2p^1$ and He $3d^3-2p^3$) or are optically thick and attenuated in the axial spectrum (Li $4p-2s$, Li $4d-2p$, He $3d-2p$).

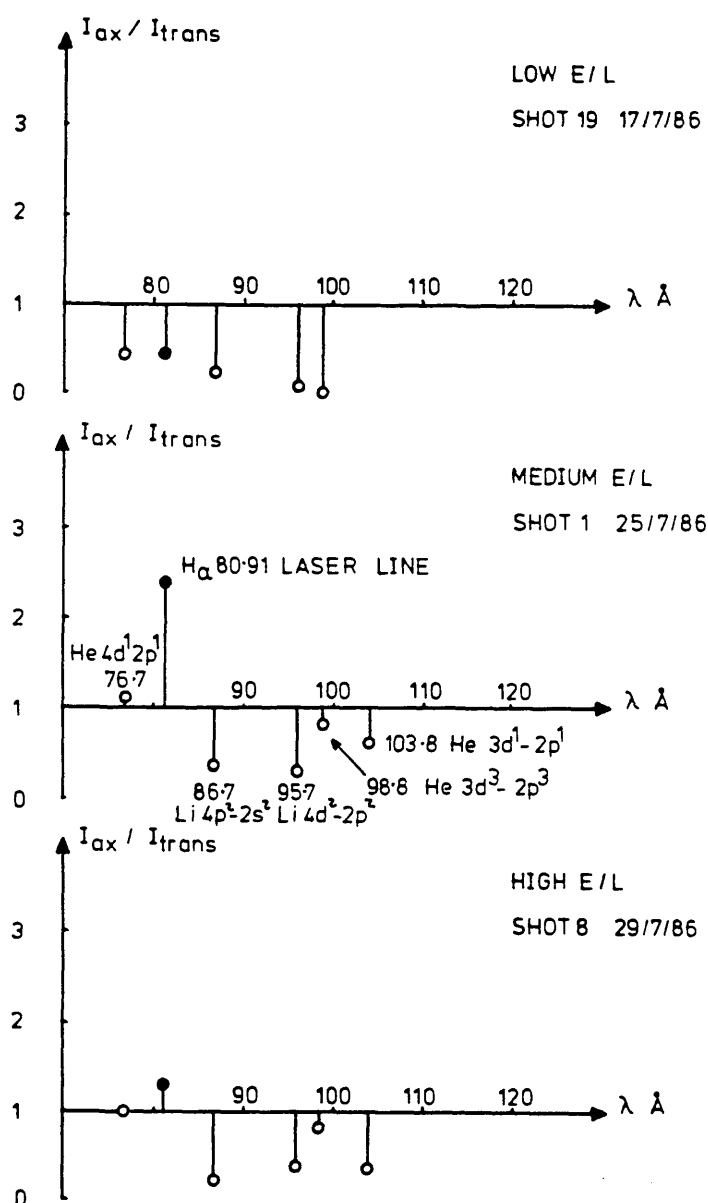


Figure 5.17 Normalised axial to transverse ratios for a range of lines, for shots in the low, medium and high E/L regimes.

5.5 COMPUTER SIMULATIONS

Comprehensive numerical simulations of the experiment have been carried out [74,75]. The method involved the improvements to the model included in the work on carbon fibres reported in chapter 4, with the exception of the Stark broadening for which no data comparable to that for carbon was available. The hope is that because static-ionic Stark broadening decreases with increasing Z , the motional Stark effect, which makes a dominant contribution to the total Stark width for carbon, will be less important for the case of fluorine.

Figure 5.18 by S. Rose [79] compares the measured and calculated time evolution of gain for a single 5.5 mm long shot from the medium E/L regime at an absorbed energy of $\approx 10 \text{ J cm}^{-1}$. For the calculation without line trapping the two are in good agreement on the magnitude and timing of peak gain, but the inclusion of line trapping significantly reduces the agreement. This gives further weight to the view that the complex radiation transport in an expanding plasma with density, temperature and velocity gradients is not well modelled with the local description of radiation trapping used in the code. The code is however in much closer agreement with experiment than previous work on carbon [8]. This is due to both additions to the physics content of the code and to improved experimental techniques which do not make assumptions about the time evolution of the source function.

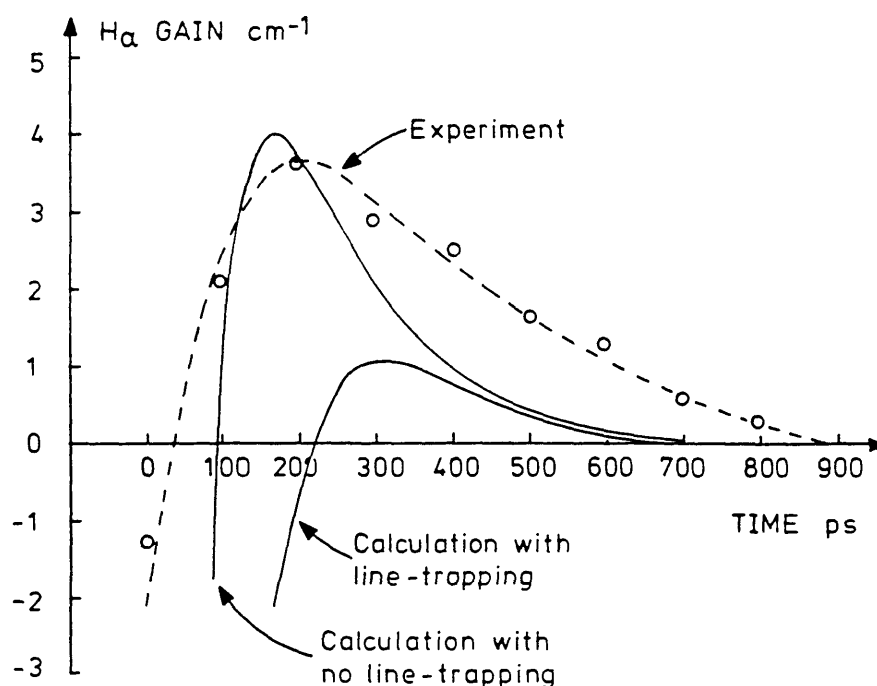


Figure 5.18 Experimental and calculated time evolution of gain for a single 5.5 mm shot from the medium E/L, high gain regime.

The data from the 1D space resolving spectrometer is both time integrated and integrated along cords through the plasma. For these reasons it is difficult to compare in detail with the radial gain profiles generated by the codes. However at early times (250 – 400 ps) the codes show peak gain at $\approx 100 \mu\text{m}$ from the fibre and this is indeed consistent with the formation of a population inversion seen in figure 5.16. Peak gain is not at the point of the largest inversion, but roughly at the highest density region where there is an inversion.

5.6 CONCLUSION

The hydrogenic recombination X-ray laser scheme was successfully scaled from carbon to fluorine. Amplification at $81 \text{ \AA}$ on the fluorine Balmer H_α line was unambiguously demonstrated using improved diagnostic techniques. The effect of coupled energy variation on the F IX laser scheme was studied. Three gain regimes were found as the plasma energy per unit length varied. These were characterised by absorption ($-1.8 \pm 0.2 \text{ cm}^{-1}$), "high gain" ($4.4 \pm 1.0 \text{ cm}^{-1}$) and "low gain" ($2.0 \pm 1.0 \text{ cm}^{-1}$) at energies of $5 \pm 0.7 \text{ J cm}^{-1}$, $9 \pm 2 \text{ J cm}^{-1}$ and $16 \pm 1 \text{ J cm}^{-1}$ respectively.

The three gain regimes were predicted by improved numerical modeling which also gave a good fit to the experimental time evolution of gain when Ly_α trapping was ignored. The hydrogenic fluorine recombination X-ray laser was the first scheme to demonstrate gain below $100 \text{ \AA}$, and is presently the best characterised scheme to have shown gain experimentally. A paper by the author [9] reporting on the work outlined in this chapter has recently been submitted to Optics Communications.

CHAPTER 6. LITHIUM LIKE ALUMINIUM X-RAY LASER EXPERIMENT

ABSTRACT

Lithium like ions are possible XUV amplifiers, and a recombination laser based on lithium like lasant ions would have significant energetic advantages over hydrogenic schemes. Gains on the order of 1 cm^{-1} at $154.6 \text{ \AA}$ have been seen by previous workers in laser produced plasmas from large aluminium slab targets. However plasmas from slab targets are believed to be subject to significant trapping of the radiative transitions which depopulate the lower laser levels. Fibre based targets can reduce trapping by enhancing Doppler decoupling in the direction transverse to the fibre axis, and thus increase the gain.

Aluminium coated $7 \text{ }\mu\text{m}$ fibre targets up to 11 mm in length were irradiated in line focus geometry with 120 ps pulses of green ($\lambda = 0.53 \text{ }\mu\text{m}$) light from the VULCAN laser at RAL. Simultaneous time resolved axial and transverse XUV spectroscopy was used to diagnose the freely expanding, adiabatically cooled Li like Al plasmas produced. At absorbed energies in the range of $4.7 \pm 1.4 \text{ J cm}^{-1}$, gains of 1.45 ± 0.5 , 1.63 ± 0.5 , 1.38 ± 0.5 , and $2.16 \pm 0.5 \text{ cm}^{-1}$ were measured for the $5d-3p$, $5f-3d$, $4d-3p$ and $4f-3d$ transitions at wavelengths of 103.8 , 105.7 , 150.7 and $154.6 \text{ \AA}$ respectively. With fibre targets, the gain of $\approx 3 \text{ cm}^{-1}$ for the $154.6 \text{ \AA}$ $4f-3d$ transition established from axial to transverse ratios was some 2.5 times higher than that seen by previous investigators using slab targets.

6.1 INTRODUCTION

The large pump power requirements of present X-ray laser schemes limits their short wavelength scalability, and future coherent X-ray sources based on such schemes will probably be confined to major central facilities. Some schemes are inherently more energetically efficient than others (see figure 6.1) and new schemes with lower pump requirements are being investigated by a number of groups. Recombination X-ray lasers require significantly less pump power to run than electron collisionally pumped neon or nickel like schemes, due to the lower ionisation energy of their laser ions. This is a major reason why the 3 TW VULCAN laser can pump an XUV amplifier operating at a shorter wavelength than the Se^{24+} scheme pumped by the 100 TW NOVA laser.

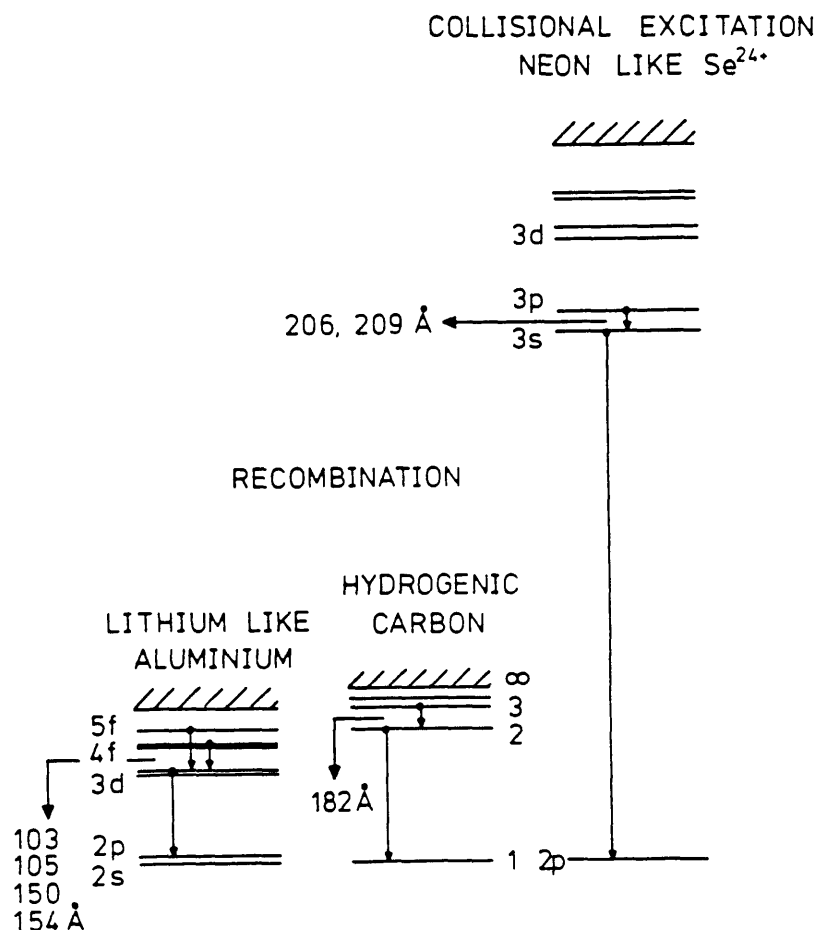


Figure 6.1 A comparison of lithium like, hydrogen like and neon like X-ray laser schemes showing the energetic advantage of using Li like ions.

From figure 6.1 it can be seen that lithium like ions are energetically an attractive proposition for use in recombination XUV lasers. Lithium like ions have a single electron in their outer shell, and as with hydrogenic ions the lithium like ion energy level spacing increases with decreasing energy. This means that in a recombination cascade, transitions between the upper levels will be collisionally dominated while transitions from lower levels will be predominantly radiative. Population inversions are thus possible in the lower levels of cooling and recombining lithium like ions.

Figure 6.2 shows a simplified term diagram for lithium like aluminium. If the 3p and 3d levels are strongly depopulated by fast radiative resonance transitions then 5-3 and 4-3 population inversions can form. As with H like schemes the thermal limit must lie above the upper laser level or collisional mixing will destroy the inversion. Similarly the lower laser level to ground decay radiation must be optically thin, or reabsorption will clog up the population inversion. If these conditions are met then 5d-3p, 5f-3d, 4d-3p and 4f-3d inversions should give gain at wavelengths of 103.8, 105.7, 150.6 and 154.6 Å respectively.

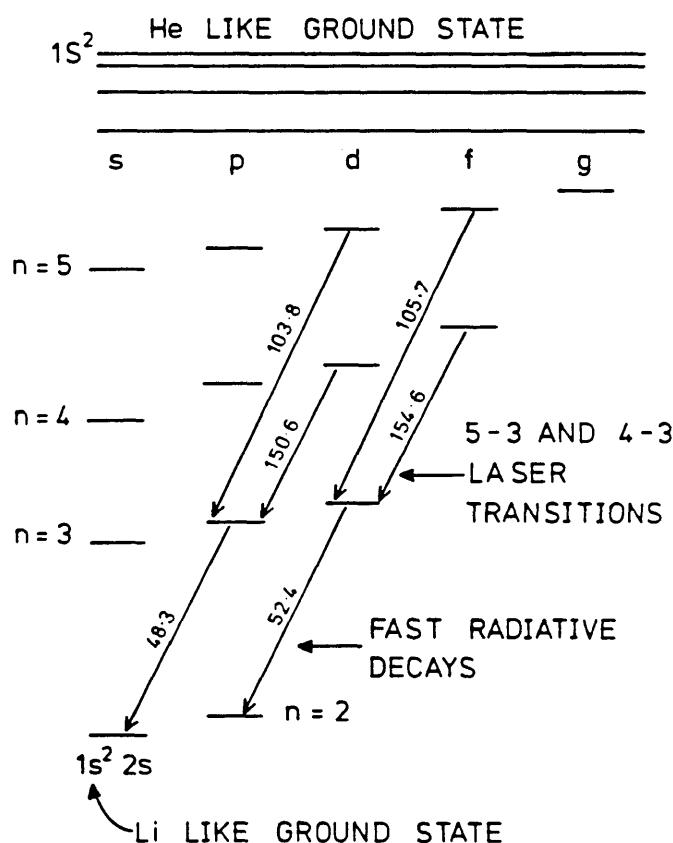


Figure 6.2 A simplified energy level diagram for Li like aluminium.

Population inversions in Li like Al plasmas produced by point focus irradiation of slab targets have been seen by a number of groups [81,82]. Jaeglé and co-workers at Palaiseau have since extended this scheme to line focus geometry at lengths up to 20 mm, and small signal gains in the order of 1 cm^{-1} in adiabatically cooled plasmas from large slab targets have been observed at $154.6 \text{ \AA}$ [39,40].

These targets were irradiated at power levels on the order 10^{10} Wcm^{-2} with relatively long (2 – 4 ns) pulses of $1.06 \text{ }\mu\text{m}$, 1ω light from a two beam Nd:glass laser system. In slab geometry a wedge shaped plasma plume expands away from the target, approximately along the path of the pump laser beam. This limits the solid angle of the blow off, and consequently decreases the expansion rate. This and the large plasma mass inherent in the slab target technique result in increased line trapping of the transitions which depopulate the lower laser levels, which in turn lowers the gain. It was suggested that a fibre based Li like Al laser would have a lower opacity due to enhanced Doppler decoupling in the radial direction, a lower plasma mass, and lower initial ground state population. Fibre targets would therefore increase the gain which could be achieved in the lithium like aluminium scheme.

A collaborative experiment involving members of Jaeglés group and the UK X-ray laser consortium was carried out to investigate XUV laser targets consisting of $7 \text{ }\mu\text{m}$ diameter carbon fibres coated with $0.35 \text{ }\mu\text{m}$ of Al. The experiment was conducted at RAL using the VULCAN glass laser and the TAE off axis line focus facility which optimises pump laser coupling into thin fibre targets. Simulations of Al fibre targets identified a coupled energy of $\approx 5 \text{ J cm}^{-1}$ and a pump laser pulse length of 120 ps for peak gain. With a coupling efficiency of $\approx 10 \%$ this energy density can be achieved with one pair of opposing beams in the TAE line focus. This allowed target lengths up to 14 mm to be investigated by spacing two 7 mm focus beam pairs along the target. However as with carbon fibres the length was restricted to less than 12 mm by bending and vibration of longer targets. For some shots less than 7 mm in length, two beam pairs run at half the usual power level were superimposed on target to improve the uniformity of irradiation.

6.2 EXPERIMENTAL

The experimental arrangement and method for this investigation was nearly identical to that used in the hydrogenic fluorine X-ray laser experiment described in chapter 4. The exceptions were the longer target length attainable (due to the lower energy requirement of this scheme) and the optimum $0.53\ \mu\text{m}$ pump laser pulse length of 120 ps identified by simulations.

The primary diagnostics were a pair of time resolved flat field XUV spectrometers used to record axial and transverse emission from the full length of the plasma column. Plasma conditions and length were observed with the usual range of plasma calorimeters, temperature monitors and pinhole cameras. Figure 6.3 shows the experimental arrangement.

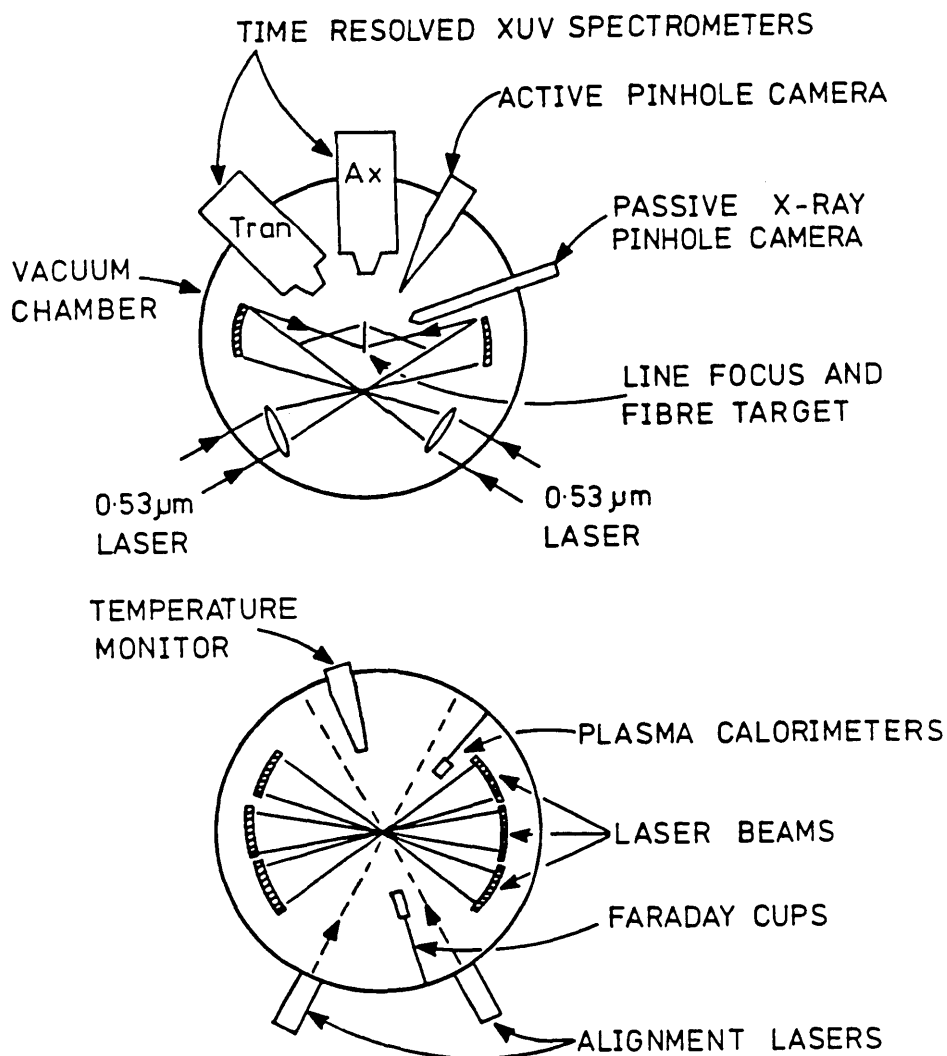


Figure 6.3 Experimental arrangement for the Li like aluminium study.

6.3 XUV SPECTROSCOPY

The pair of flat field spectrometers were run up in time integrated mode using a phosphor plate detector and image intensifier. Figure 6.4 shows a typical time integrated transverse spectrum from an Al coated glass slab target. The complexity of the spectrum compared to lower Z hydrogenic spectra is immediately obvious. Resolving the large number of lines in high Z spectra has been found to be a problem with the flat field instruments due to their relatively low ($0.5 \text{ \AA}$) resolution. The Al spectrum was identified by cross reference to well understood spectra from carbon targets.

As with time integrated spectra from CH and LiF targets, the spectrum is again dominated by short wavelength resonance transitions in multiple orders. The dominant features are multiples of the 2p-3d and 2s-3p resonance transitions at 52.4 and 48.3 $\text{\AA}$. This is encouraging as these are the two transitions which depopulate the 3d and 3p lower laser levels. The possible 3d-f4 and 3p-4d laser transitions at 154.6 and 150.6 $\text{\AA}$ are just resolved, however the 3d-5f and 3p-5d laser lines at 105.7 and 103.5 $\text{\AA}$ are obscured by the second order 2s-3p resonance line. The remainder of the spectrum is composed of many poorly resolved Be and He like aluminium lines.

Time resolved spectra were obtained by replacing the phosphor plates with streak cameras on both spectrometers. As with the fluorine experiment it was found that the transverse spectrum was significantly source broadened compared to the axial spectrum. This and the high brightness of the resonance lines in the transverse spectrum due to radial Doppler decoupling made it difficult to resolve the 5f-3d and 5d-3p laser lines. For this reason the spectrometers were set to observe longer wavelengths, and recorded these two lines in second order, increasing the separation between them and the resonance line.

Figure 6.5 compares axial and transverse streak spectra from an Al fibre shot. The four prospective laser lines are the most prominent features in the axial spectrum and are relatively long lived compared to the resonance lines which depopulate the lower laser levels. The laser lines all peak in brightness at around 300 - 500 ps after the laser pulse, somewhat earlier than was seen for hydrogenic carbon and fluorine plasmas. This may be due to a more rapid cooling of the plasma by radiation loss, which scales as Z^4 for line emission and Z^2 for Bremsstrahlung and radiative recombination.

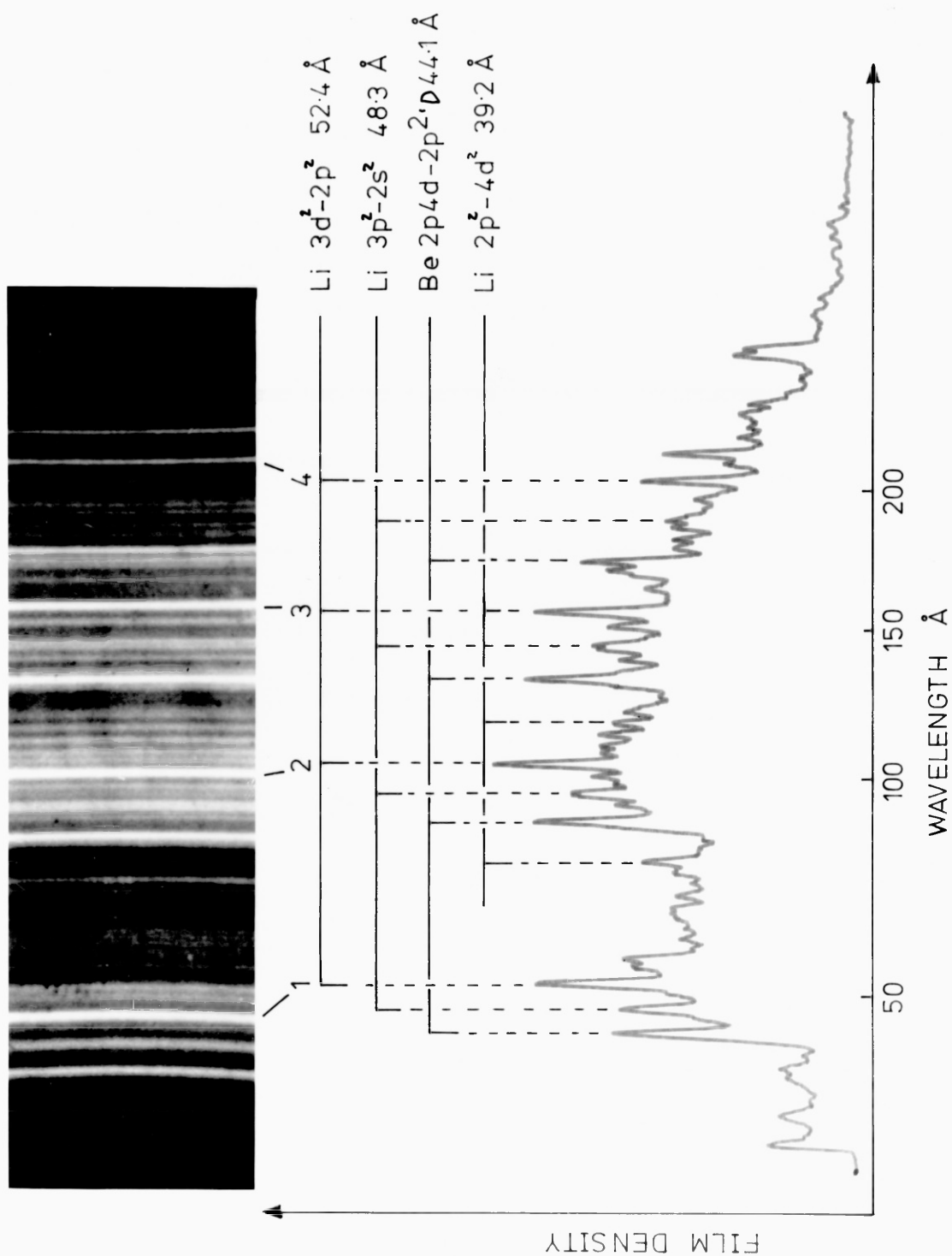


Figure 6.4 Time integrated XUV spectrum from an aluminium coated glass slab target

Li LIKE Al STREAK SPECTRA

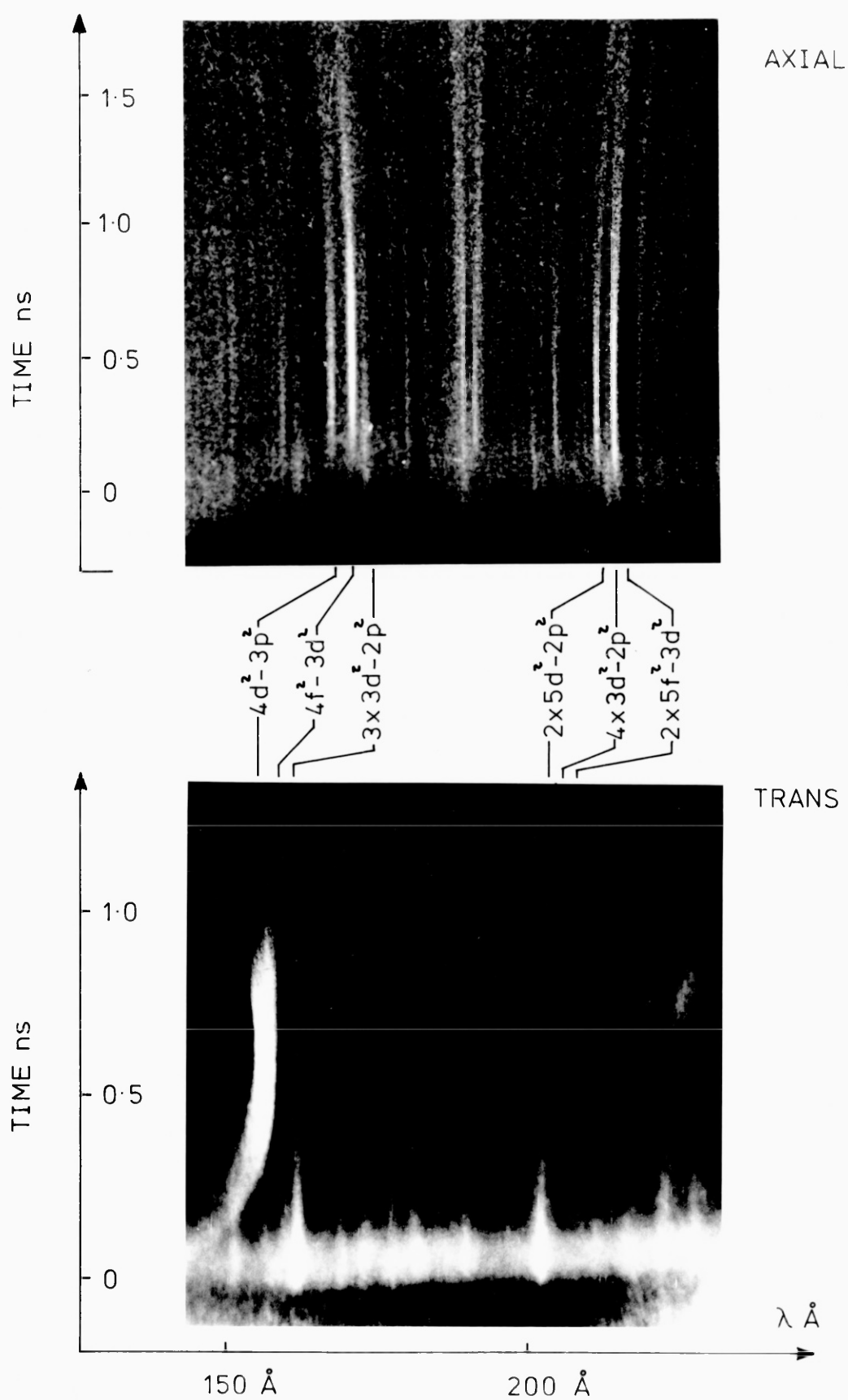


Figure 6.5 Axial and transverse XUV streak spectra from an Al coated fibre target.

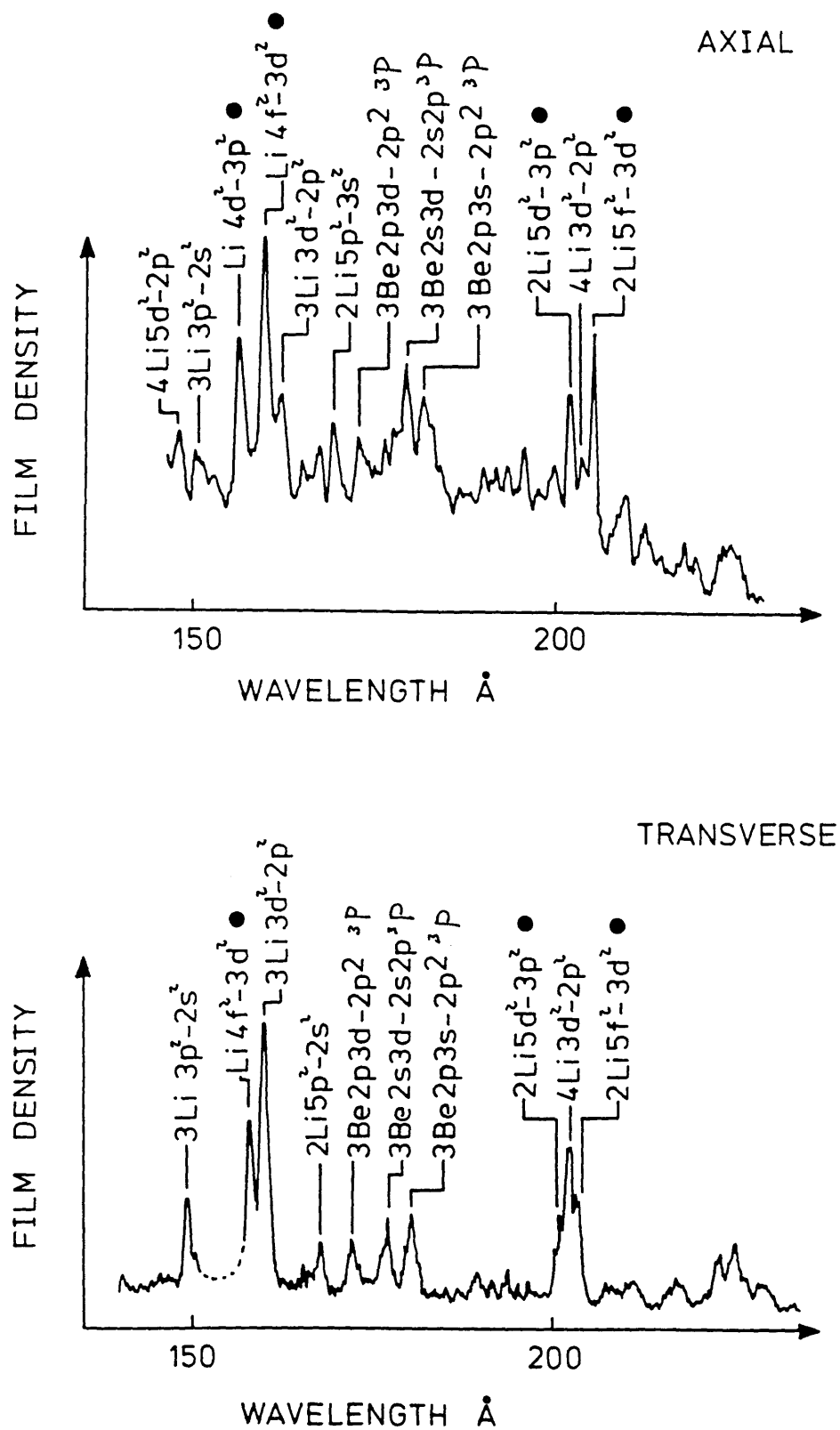


Figure 6.6 Densitometer traces of the axial and transverse streak spectra presented in figure 6.5.

In contrast, the transverse streak is dominated by the resonance transitions and the whole spectrum is short lived compared to the axial emission. As noted earlier, the high brightness of the resonance lines and the source broadening of the transverse spectrum made the prospective laser lines difficult to observe. This can be seen quite clearly in figure 6.6 which shows densitometer traces of the two streaks in figure 6.5. All four laser lines are well resolved in the axial spectrum. However in the transverse spectrum the second order 5f-3d and 5d-3p lines are only seen as shoulders on the fourth order 2p-3d resonance line. The first order 4f-3d and 3p-4d lines are somewhat better resolved. The remainder of the lines in the transverse spectrum are higher order Be like aluminium lines. No strong He like lines were seen in the streak spectra indicating that the plasma was not overstripped during the heating pulse. Similarly no carbon lines were seen in either axial or transverse spectra showing that the 0.35 μm overcoat of Al did not burn through to the fibre substrate.

6.4 SELECTION OF THE DATA SET

The aluminium coated fibre targets were found to be difficult to align in the TAE line focus using the usual CW argon ion laser. This was due to the CW laser heating the fibres and causing them to bend out of the line focus in a similar way to a bi-metallic strip. This resulted in a wide shot to shot variation in coupled energy over the whole data set. The final data set of 11 shots used for the investigation of gain was selected on the basis of absorbed plasma energy. Only shots with coupled energies in the range $4.7 \pm 1.4 \text{ J cm}^{-1}$ which simulations identified as the regime for peak gain were considered.

In previous experiments on carbon and fluorine plasmas the absorbed energy was assigned on the basis of the plasma calorimeter readings alone. Unfortunately these devices are known to be rather inaccurate on a shot to shot basis, due to the difficulty in distinguishing between the large scattered pump laser signal and the relatively small signal from the expanding plasma. However the integrated pinhole camera brightness is also a good indication of the absorbed energy and was used in preference on this experiment. A best fit line was found to a plot of the integrated APC brightness versus the absorbed energy as given by the plasma calorimeters (see figure 6.7). Absorbed energy was then inferred from the APC brightness of each shot rather than from the individual calorimeter readings. As in previous experiments the plasma length was taken to be that from which $\approx 90 \%$ of the emission was seen in the APC images.

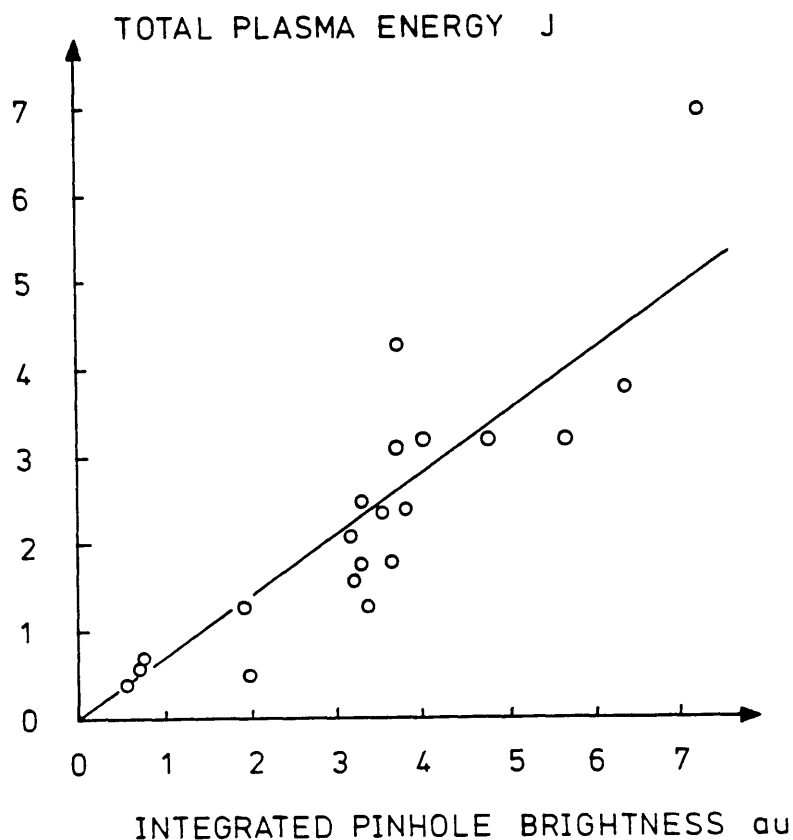


Figure 6.7 Integrated active pinhole camera brightness versus absorbed energy.

6.5 OBSERVATION OF GAIN FROM AXIAL TO TRANSVERSE RATIOS

As was seen in section 6.3 the 5d-3p, 5f-3d and 4d-3p transitions were poorly resolved in the transverse spectra, making an axial to transverse comparison for these lines difficult. However the 4f-3d line at $154.6 \text{ \AA}$ was far enough away from the resonance lines to be well resolved in both the axial and transverse spectra from fibre targets. To make a comparison between axial and transverse spectra the two spectrometers must be cross calibrated. As in chapters 4 and 5 this was attempted using an angled slab target irradiated by a very short (1 mm) section of the line focus. However the large source broadening inherent in plasmas from this type of target made it impossible to resolve any of the possible laser transitions and the spectrometers could not be cross calibrated in this way.

An alternative method due to Regan [83] was used. For a range of times the uncalibrated axial to transverse ratio of the 154.6 Å line was plotted against plasma length, using a series of six shots with similar plasma energies per unit length. From equation (4.3) it follows that the uncalibrated axial to transverse ratio is simply :-

$$\frac{I_{ax}}{I_{trans}} = \frac{e^{\alpha \ell} - 1}{k \alpha \ell} \quad (6.1)$$

$$k = \frac{\text{axial sensitivity}}{\text{transverse sensitivity}} \quad (6.2)$$

As can be seen in figure 6.8, at 200 ps after the heating pulse the axial to transverse ratio is nearly constant for a range of lengths, and equation (6.1) indicates an optically thin line with no gain. The intercept at $\ell = 0$ gives the value of k at 154.6 Å. However at later times the axial to transverse ratio can be seen to increase exponentially with plasma length. This is a conclusive indication of gain, and numerical fitting of the points in figure 6.8 to equation (6.1) gives a gain coefficient of $2.5 \pm 0.5 \text{ cm}^{-1}$ at 400 ps after the laser pulse. A similar analysis of the other prospective laser lines was impossible due to the masking of these lines by resonance transition in the transverse spectra.

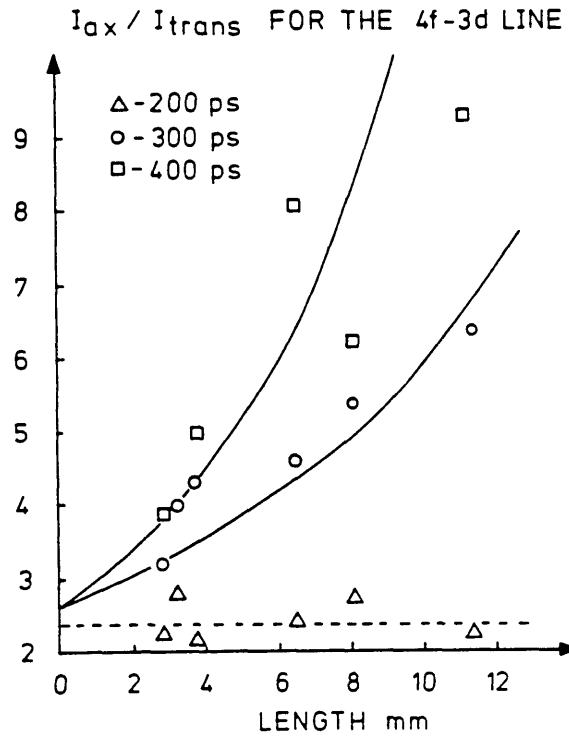


Figure 6.8 Axial to transverse ratio for the 154.6 Å line at several times

6.6 TIME EVOLUTION OF THE 4f-3d GAIN FROM AXIAL TO TRANSVERSE RATIOS

The rapid decay of transverse emission in time made it impossible to follow the time evolution of gain for the 4f-3d transition after 400 ps. However the large opacity for this line at early time when the plasma is too hot and dense to support gain was well observed. Figure 6.9 shows the 4f-3d gain as a function of time found for three different length shots by solving equation (6.1) numerically. In all cases gain turns on sharply at ≈ 200 ps after the peak of the laser pulse. Gains on the order of $2.5 - 3 \text{ cm}^{-1}$ are seen on all these shots at ≈ 400 ps, the latest time that was observable using this technique.

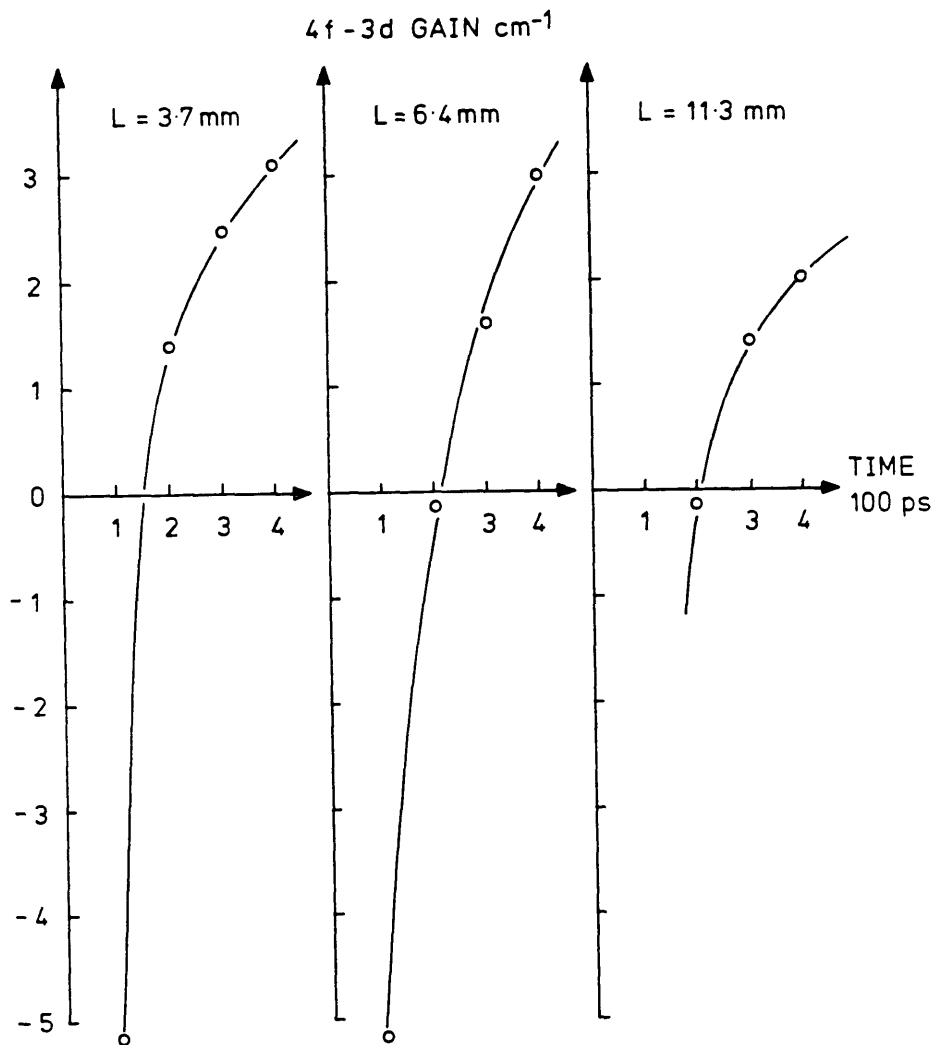


Figure 6.9 Time evolution of gain for the 4f-3d transition found from axial / transverse ratios for three different length shots.

A slight decrease in the gain found from axial to transverse ratios with increasing plasma length can also be seen in figure 6.9. This was a general trend seen on many shots, and is believed to be due to the plasma in short length shots expanding more nearly spherically than cylindrically. This results in a faster expansion of the plasma, faster cooling, and consequently higher gain at early times in short length shots.

6.7 OBSERVATION OF GAIN FROM AXIAL INTENSITY VARIATION

As in chapters 4 and 5 the gain can also be found from the axial variation of intensity with plasma length. Equation (2.33) shows that below saturation, the axial intensity is given by $I_{ax} = S(e^{\alpha\ell} - 1)$ where S is the plasma source function.

Figures 6.10 – 6.13 are plots of axial brightness against plasma length for the four prospective laser lines at the times of peak gain. These times were identified by similar plots over a range of times (section 6.8). In all cases exponential growth can clearly be seen, and the gains for these transitions were found by fitting the experimental growth curves to an equation of the form $I = S(e^{\alpha\ell} - 1)$.

In contrast figure 6.14 shows a plot of axial brightness versus plasma length for the 3rd order 52.4 Å 3d–2p and 48.3 Å 3p–2s resonance lines. No gain is expected on these lines, and figure 6.14 shows evidence of saturation, as would be expected for lines optically thick in the axial direction.

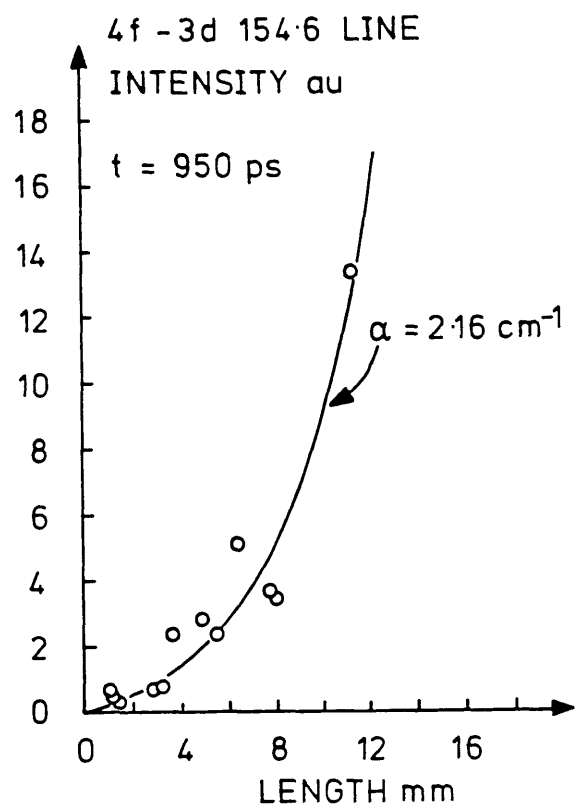


Figure 6.10 4f-3d axial brightness versus length showing exponential growth.

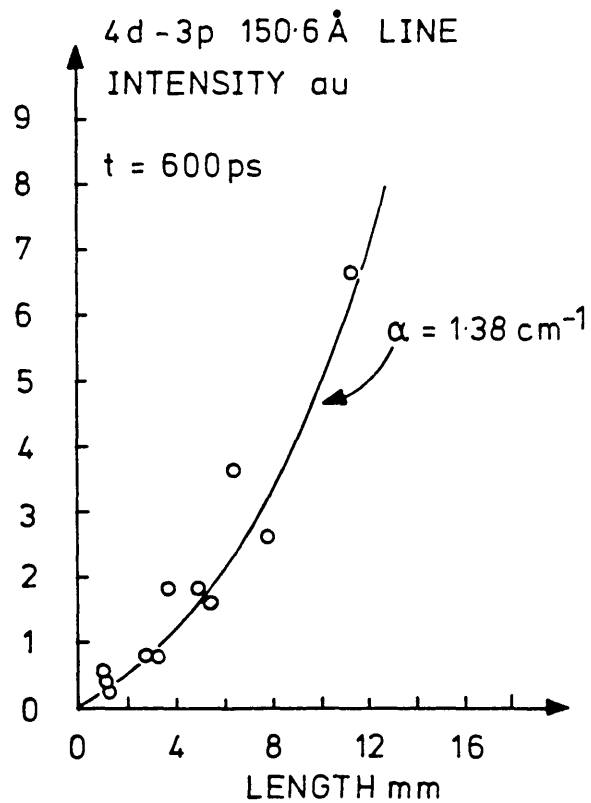


Figure 6.11 4d-3p axial brightness versus length showing exponential growth.

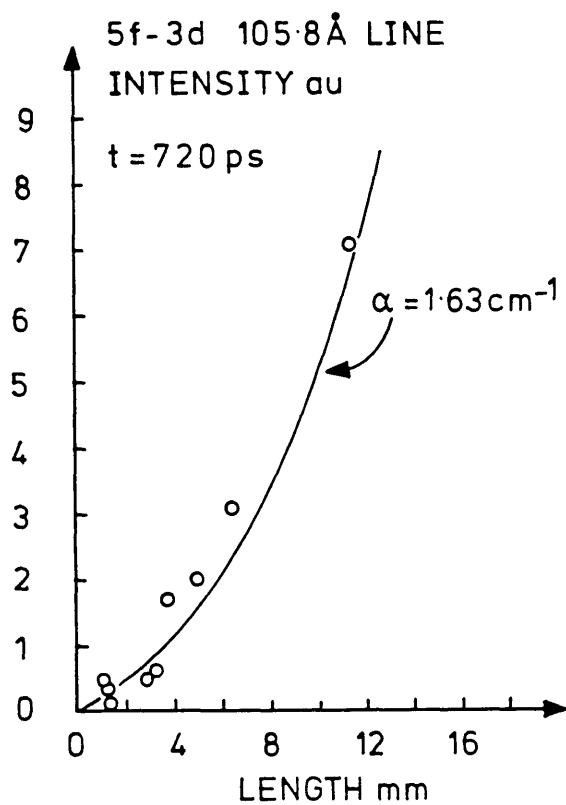


Figure 6.12 5f-3d axial brightness versus length showing exponential growth.

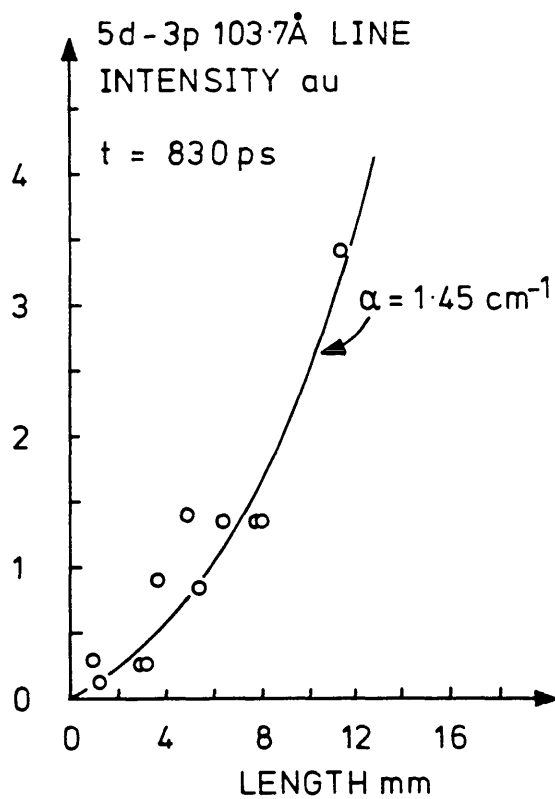


Figure 6.13 5d-3p axial brightness versus length showing exponential growth.

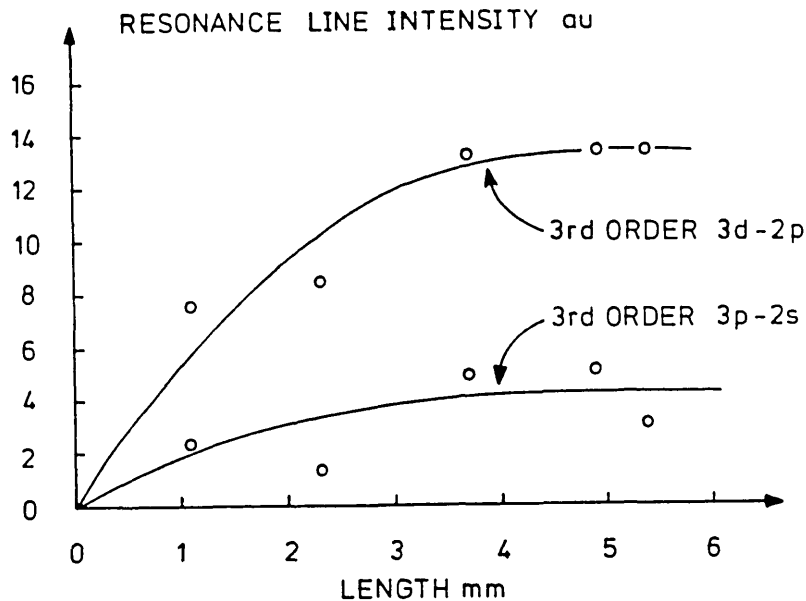


Figure 6.14 3rd order 3d-2p and 3p-2s resonance lines showing saturation with increasing plasma length.

6.8 TIME EVOLUTION OF GAIN FROM AXIAL SPECTRA

Axial to transverse ratios could not be used to follow the time evolution of gain for the four lithium like aluminium laser lines as was done for the hydrogenic fluorine and carbon schemes. The temporal history of gain can also be found by comparing the axial brightnesses of long and short shots with similar plasma conditions. However in chapter 4 it was shown that this method does not treat the time evolution of the source function or the expansion of short length plasmas correctly, and this can result in significant errors.

A second method for analysing the axial data which is less subject to these errors was used by Regan [83] to find the gain of these four laser lines as a function of time. Curve fittings of the axial data to growth equations of the form $I(t) = S(t)(e^{\alpha(t)l} - 1)$ were carried out for the full data set rather than just two shots over a range of times. Figures 4.15-4.18 show the time evolution of gain for the laser lines found using this method.

All four laser transitions clearly show gain, with the gain turning on ($\alpha > 0$) at around 200 ps after the peak of the heating pulse. The 5d-3p, 5f-3d and 4f-3p transitions all peak at 450 - 500 ps after the heating pulse, and the 4f-3d transition significantly later at 800 ps. The axial emission for these lines peaks at ≈ 300 ps, which is well before the time of maximum gain. This is the reverse of the hydrogenic case, where in both the carbon and fluorine experiments peak gain occurs before the peak of the axial emission.

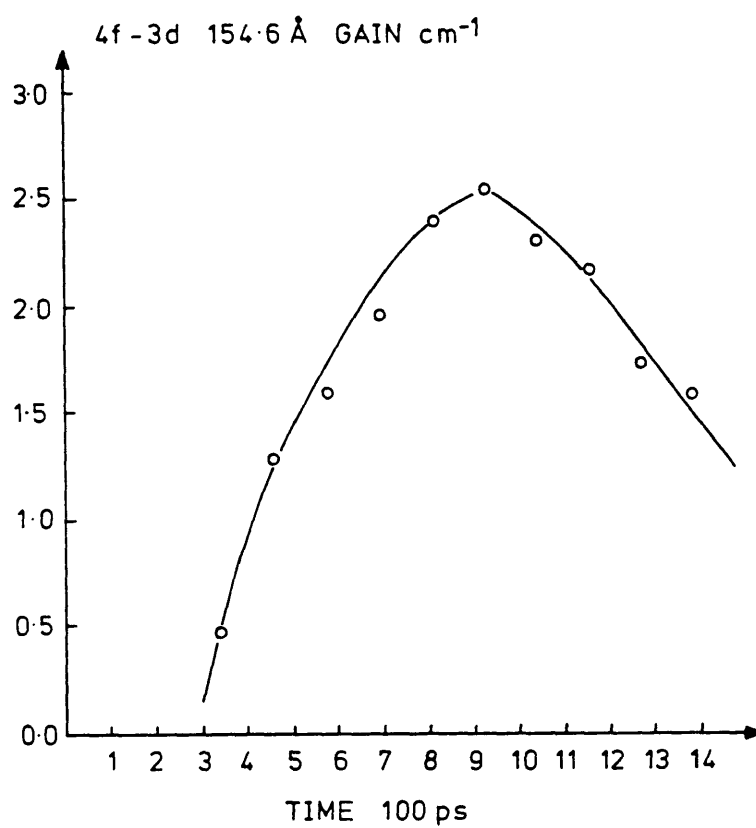


Figure 6.15 4f-3d time evolution of gain found from axial spectra.

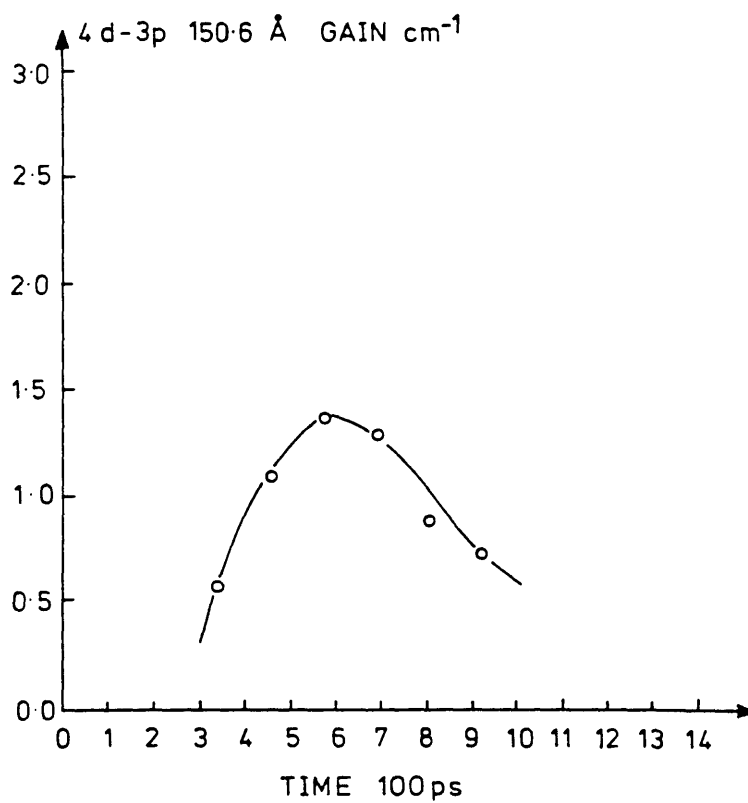


Figure 6.16 4d-3p time evolution of gain found from axial spectra.

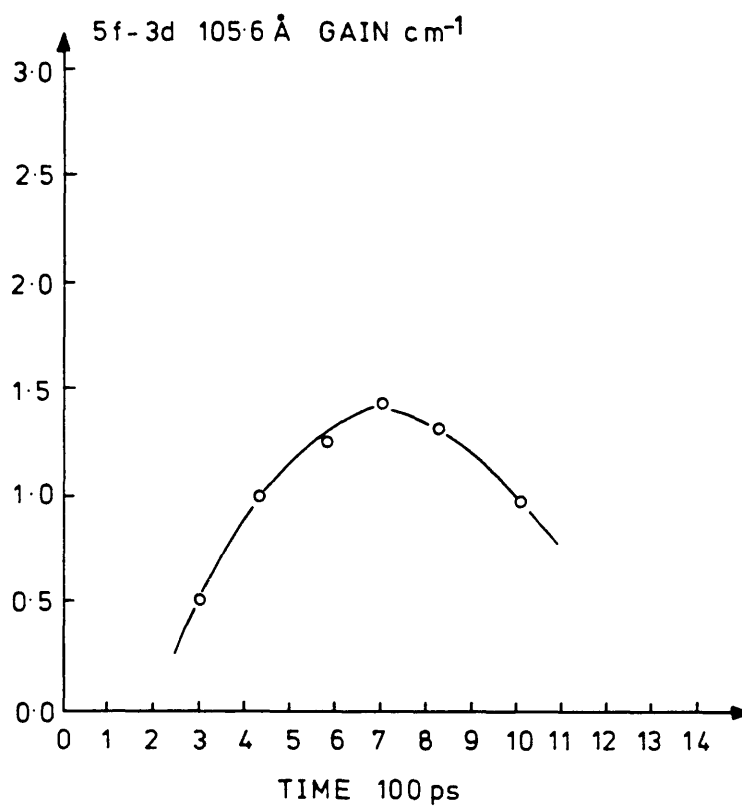


Figure 6.17 5f-3d time evolution of gain found from axial spectra.

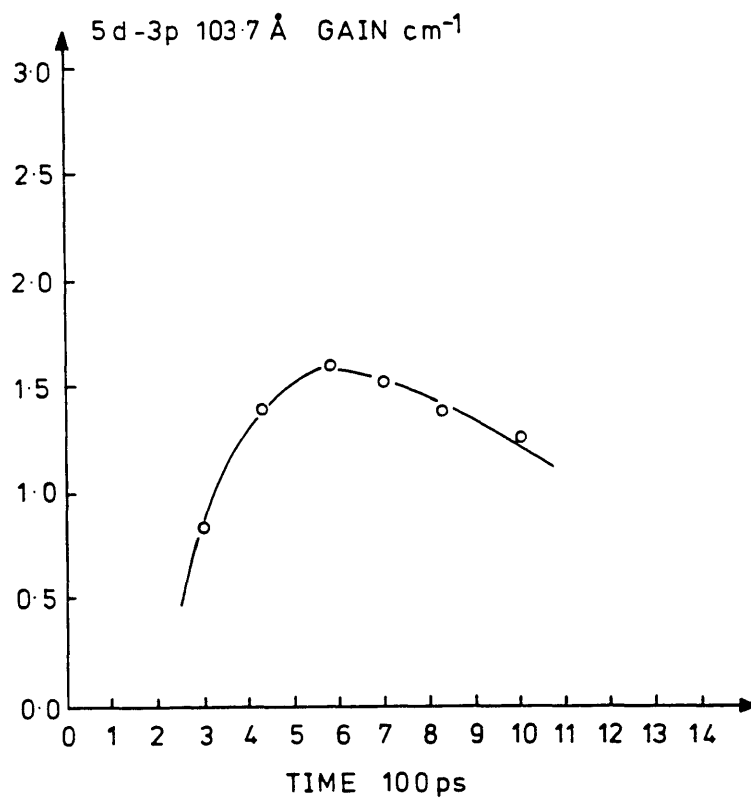


Figure 6.18 5d-3p time evolution of gain found from axial spectra.

6.9 COMPUTER SIMULATIONS

The simulation of Li like X-ray laser schemes is at a less advanced stage than that for H like schemes due to the greater complexity of the atomic physics. Initial calculations which only included Li like levels up to $n = 5$, with l but not j splitting indicated a large ($\approx 11 \text{ cm}^{-1}$) gain for the $4f-3d$ line but very small (0.1 cm^{-1}) gain for the $5f-3d$ line. No gains were predicted for the $d-p$ transitions using this code.

The atomic physics model has since been considerably extended by Pert and Henshaw [84]. The He like ion is now represented by four levels, and the Be like ion by a further eight. The Li like ion is now described by 37 levels, with full l and j splitting up to $n = 5$, l splitting for $n = 6$, and H like states for $n = 7 - 41$. This code predicts gains of $\approx 3 \text{ cm}^{-1}$ for $4f-3d$, 0.9 cm^{-1} for $5f-3d$, 1.8 cm^{-1} for $4d-3p$ and 0.4 cm^{-1} for $5d-3p$. The 4-3 gains are in fair agreement with experiment as can be seen in figure 6.19 which compares the calculated time evolution of gain for the $4f-3d$ line at $154.6 \text{ \AA}$ with that found experimentally using axial and axial to transverse ratio methods.

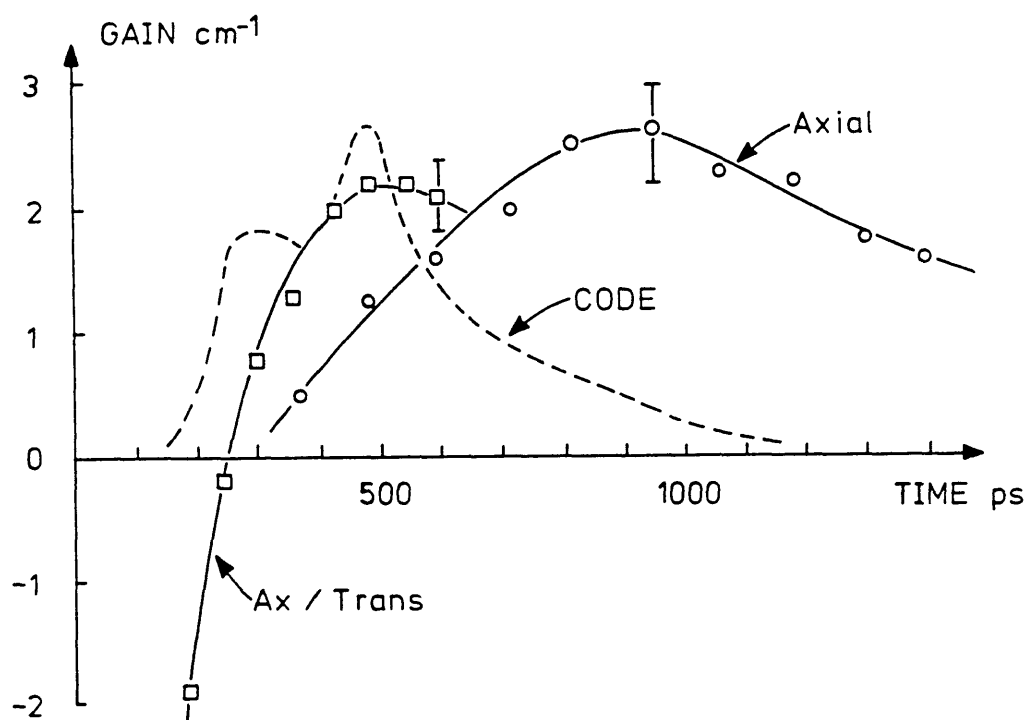


Figure 6.19 Calculated and measured time evolution of gain for the $4f-3d$ transition at $154.6 \text{ \AA}$. Measured gains from both axial only and axial to transverse ratio methods.

6.10 CONCLUSION

The lithium like aluminium XUV laser scheme was investigated experimentally, using fibre targets for the first time. Simultaneous axial and transverse time resolved spectroscopy was used to conclusively demonstrate gain on the order of 3 cm^{-1} at $154.6 \text{ \AA}$ on the $4f-3d$ transition of Li like Al. This was some 2.5 times higher than the gain seen in laser produced Al plasmas from massive slab targets. This increased gain is believed to be due to the lower mass, lower ground state population and enhanced Doppler decoupling of the lower state decay radiation inherent in plasmas from thin fibre targets.

Using only axial time resolved spectroscopy smaller gains of 1.45 ± 0.5 , 1.62 ± 0.5 , $1.38 \pm 0.5 \text{ cm}^{-1}$ and $2.16 \pm 0.5 \text{ cm}^{-1}$ were conclusively demonstrated for the $5d-3p$, $5f-3d$, $4d-3p$ and $4f-3d$ transitions of Li like Al. The exponential increase of axial intensity with plasma length was clearly observed for these lines.

The time evolution of gain for all the laser lines was found from curve fitting to the full axial data set rather than from pairs of shots. This is believed to give a more accurate picture of the gain history than the long / short shot comparison since this method is less subject to errors from plasmas with different source functions. In addition, the time evolution of gain for the $4f-3d$ line was found by comparing axial to transverse ratios on single shots. As with experiments on carbon and fluorine, this technique was found to give the best fit to code simulations.

In contrast to the hydrogenic recombination schemes, peak gain was seen to occur after the peak of the axial emission of the laser lines. The more rapid cooling of aluminium plasmas compared to carbon and fluorine was thought to be due to the strong Z scaling of radiation loss.

CHAPTER 7. RADIATION COOLING FOR THE HYDROGENIC FLUORINE LASER SCHEME.

ABSTRACT

It has been suggested that small amounts of high Z dopants could significantly increase the radiation loss from adiabatically cooled recombining laser plasmas [48,85,86]. In an XUV amplifier this would result in a more rapid cooling of the plasma and population inversion would form at a higher density, enhancing both the gain and the short wavelength scalability of such schemes.

Radiation cooling for the hydrogenic fluorine "laser" was investigated experimentally by the author. Slab targets of various fluoride salts were irradiated in point focus geometry with 2 ns pulses of $0.25\ \mu\text{m}$ KrF light from the SPRITE laser. Time resolved X-ray and XUV spectroscopy was used to diagnose the plasmas produced. The time evolution of the plasma electron temperature was found from the ratio of F Ly_γ to F He_γ and by comparison of experimental spectra to synthetic spectra generated by an atomic physics code. Hydrocode simulations were also used to model the plasma heating phase, and good agreement with the experimental data was found.

For NaF and MgF_2 where the atomic number of the dopant was similar to fluorine ($Z = 9$), cooling rates similar to that for LiF targets were measured. However for KF and CaF_2 targets where the dopant has a significantly higher atomic number, a factor of 2 – 3 enhancement in the cooling rate was seen.

7.1 INTRODUCTION

With adiabatic expansion as the sole cooling mechanism, the strong scaling of initial density with atomic number places a serious limitation on the shortest wavelength at which a fibre based hydrogenic recombination laser could be operated. If the cooling of the plasma gain medium could be enhanced by radiation loss then the initial density requirement would be relaxed somewhat. The hydrogenic scheme could then be more easily scaled into the water window (H like Al $H_{\alpha} = 39 \text{ \AA}$). At wavelengths longer than this the gain would be enhanced as a population inversion would form in a higher density plasma.

Radiation loss by Bremsstrahlung and radiative recombination scales as Z^2 (Z the atomic number of the radiating species), while line radiation loss scales as Z^4 . Radiation loss has been identified by a number of groups [48,85,86] as an important mechanism for cooling recombining laser produced plasmas. It has been suggested that small amounts of high Z dopants would significantly increase the radiative cooling in laser pumped XUV amplifiers. In addition the dopant would increase the electron density and hence the recombination rate without adversely affecting the opacity of the transitions which depopulate the lower laser level.

Seely and co-workers at NRL [48] have reported a two fold enhancement of gain at $182 \text{ \AA}$ in laser produced carbon plasmas from thin CH foils coated with $2000 \text{ \AA}$ of selenium ($Z = 32$), as compared to plain CH foils. Both radiative cooling by Ne like selenium ions and a higher electron density were thought to have contributed to this increase in gain. Suckewer's group at Princeton [35] have operated a $182 \text{ \AA}$ hydrogenic carbon "laser" in which the primary cooling mechanism was radiation loss. Carbon plasmas produced by long pulse CO_2 laser irradiation of disc targets were confined by a magnetic field preventing radial expansion, and gains on the order of 5 cm^{-1} were observed.

These experiments have shown that radiation cooling can have a significant effect on the operation of hydrogenic ASE amplifiers. The possibility of enhancing the gain of the hydrogenic fluorine laser by doping targets containing fluorine with higher Z elements was investigated experimentally by the author. High Z elements could be incorporated into X-ray laser targets by ion implantation, vapour deposition of thin layers and so on. However for simplicity a range of fluoride salts compressed into solid slabs were investigated.

These targets were irradiated in point focus geometry with 2 ns long pulses of 0.25 μm KrF light from the SPRITE laser system. Time resolved X-ray and XUV spectroscopy was used to study the ionisation balance and cooling rates of the plasmas produced, and to search for any line overlaps which might repopulate the $n = 2$ lower laser level of the F IX ion.

The time evolution of the electron temperature was found by comparing experimentally measured F Ly $_{\gamma}$ to He $_{\gamma}$ line ratios with those predicted by the atomic physics code RATION [87] used in conjunction with the spectral modelling code SPECTRA [87]. These two codes were also used to estimate the effect of opacity on the measured line ratios. The 1D hydrocode MEDUSA [89] was used to model the laser heating phase of the experiment, and to provide estimates of electron density and temperature profiles for inclusion in the atomic physics models.

The original version of MEDUSA does not include radiation transport, and was not used to model the cooling phase of the plasma after the turn off of the laser pulse. A newer version of MEDUSA which does include radiation, developed by J. Edwards [90], is presently being used to model the enhancement of cooling by high Z dopants observed in this experiment.

7.2 EXPERIMENTAL

Ideally this experiment should have been carried out with a heating pulse on the order of 70 ps, as was used in the fibre based X-ray laser experiments reported in previous chapters. However the SPRITE KrF laser system was configured to produce ≈ 2 ns pulses at the time of the experiment. This is the same sort of timescale as that over which we wish to study the cooling of laser produced plasmas, and this made it somewhat difficult to observe any enhancement in cooling rates due to high Z dopants.

An F/10 aspheric doublet lens was used to focus 2 ns, 20 J pulses of 0.25 μm KrF laser light from SPRITE onto slab targets of various fluoride salts. Targets were made by compressing LiF, NaF, MgF $_2$, KF, CaF $_2$ and SrF $_2$ into thin wafers. An 80 μm tungsten wire was used as a surrogate target on which to focus the lens. The focal point was found by using the tip of the wire to obscure a reference KrF laser beam which was colinear with the main SPRITE beam.

The position of the wire was then referenced with two orthogonal HeNe laser beams, and the wire replaced by a slab target. The alignment procedure was carried out under vacuum to avoid any problems with refraction by air in the target chamber. The 80 μm diameter focal spot resulted in a peak irradiance on target of $\approx 1.5 \times 10^{14} \text{ Wcm}^{-2}$. The size of the focal spot was measured with active and passive pinhole cameras, and laser beam calorimetry was used to record the energy incident on target.

X-ray emission from the laser produced plasma was recorded over a 10 to 40 $\text{\AA}$ spectral range with a flat field XUV / X-ray spectrometer. A 2400 lines per mm Hitachi grating set at 1.5° grazing incidence was used to disperse the spectrum, and gave a resolution of $\approx 300 \text{ m\AA}$ over this range. The spectrum was time resolved by a low-magnification ($\times 0.8$) streak camera with a low density CsI photocathode. The camera streak rate of $\approx 280 \text{ ps / mm}$ corresponded to 225 ps / mm on film after the output of the camera had been demagnified by an image intensifier. With a 0.5 mm timing slit, this gave the instrument a temporal resolution of $\approx 110 \text{ ps}$ over a 7 ns window.

The spectrometer was protected from plasma blow off and reflected laser light (to which the CsI photocathode is very sensitive) by a 2.0 μm thick Al foil filter. The filter transmittance varies nearly linearly from 20 % to 60 % with wavelength over the spectral range 18 – 11 $\text{\AA}$ in which the H like and He like fluorine resonance lines of interest lie.

A second time resolved XUV spectrometer which used a 1200 lines / mm flat field grating at 3° grazing incidence recorded emissions in the 50 – 150 $\text{\AA}$ range. This instrument provided little useful information as the thickness of Al foil filtering required to shield the streak camera from stray laser light also resulted in serious attenuation of XUV radiation over the spectral range of the instrument. Both time resolved spectrometers were set up to observe the slab target at $\approx 45^\circ$ to the incident laser beam. Figure 7.1 shows a schematic of the experimental arrangement.

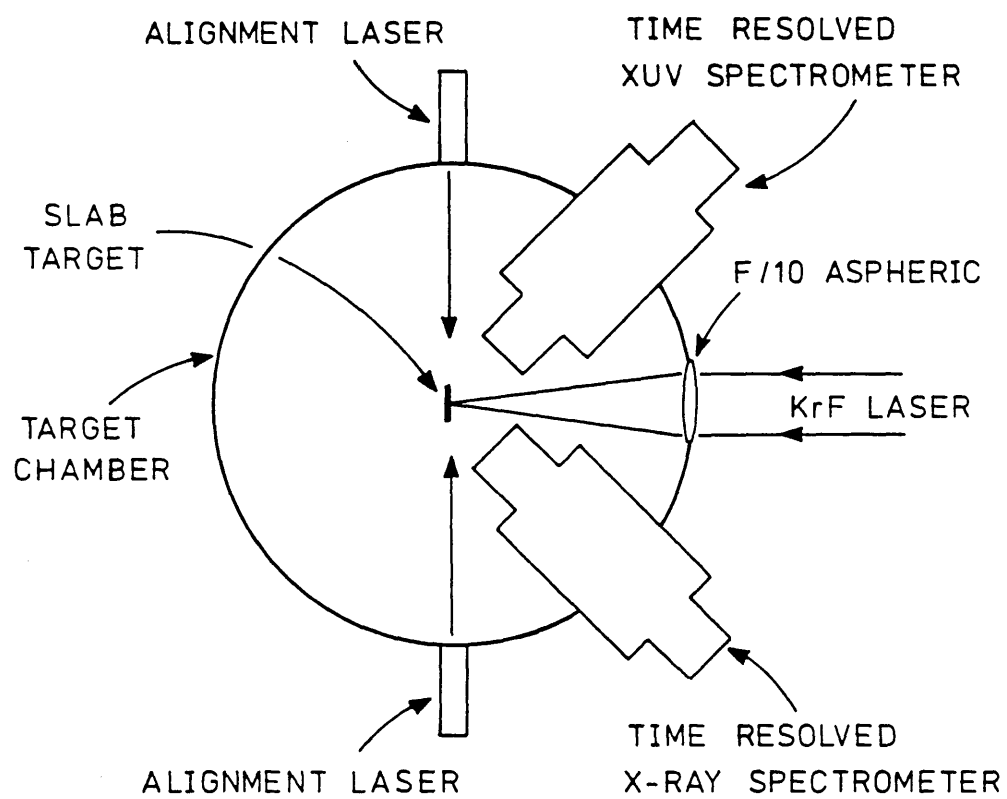


Figure 7.1 Diagnostic arrangement for the radiative cooling study.

7.3 X-RAY SPECTROSCOPY

Time resolved X-ray streak spectra from LiF, NaF, MgF₂, KF, CaF₂, and SrF₂ slab targets were recorded with the 2400 lines / mm spectrometer. This was the first such use of this particular instrument to diagnose laser produced plasmas, and has recently been reported on by Kiehn et-al [91]. Due to the time limits on this experiment, only a few shots were recorded for each target material.

Figures 7.2 to 7.7 show X-ray streak spectra and densitometer traces from the full range of targets investigated. In all cases the cut off on the short wavelength side is due to the position of the streak camera entrance slit in the flat field of the spectrometer, while the cut off to long wavelength is believed to be due to a combination of decreasing grating efficiency and rapidly increasing attenuation by the Al foil filter. The turn on of the X-ray lines in time is rather indistinct compared to streak spectra recorded on shorter pulse X-ray laser experiments, and this is due to the long rise time of the 2 ns SPRITE pulse.

The spectrum from the LiF target (figure 7.2) clearly shows the H like fluorine Lyman series and He like resonance lines. The increase in both background and line intensity to shorter wavelength is believed to be due to decreasing attenuation by the Al filter. The spectra from NaF and MgF₂ targets (figures 7.3 and 7.4) are again dominated by H and He like fluorine lines. In figure 7.3 the first members of the sodium Ly and He series can also be seen, but the remainder of these series and the magnesium Ly and He series members all lie beyond the short wavelength cut off of the spectrometer. It is not surprising that the spectra from NaF and MgF₂ targets are rather similar to each other and to the LiF spectrum as the atomic numbers of sodium (11) and magnesium (12) are rather close to that of fluorine (9).

In contrast the spectra from KF (Z = 19) and CaF₂ (Z = 20) targets (figures 7.5 and 7.6) have a larger background and contain significant contributions from Li and Be like states of the higher Z ions, reflecting the higher ionisation potentials of potassium and calcium. In both cases the fluorine Ly and He series are still well observed however, showing that the addition of some higher Z doping has not resulted in too cold a plasma.

The spectrum from the SrF_2 ($Z = 28$) target (figure 7.7) is dominated by a large number of unidentified (probably Ne like) Sr lines and has a very large background. F Ly_α , Ly_β and He_β are visible but the higher series members are lost in the noise. The small Ly_α and He_α line intensities also suggests that the plasma may be over cooled.

Over the range of the data set, a general trend of increasing background with increasing Z of the dopant ion is seen. For dopants of $Z \leq 20$ (calcium) the fluorine Ly and He series are well observed, and no line overlaps with F Ly_α (which would repopulate the ground state of a 2-1 inversion in H like fluorine) are seen. This is not the case for targets containing strontium, where the spectrum is dominated by a large background and many Sr lines.

The line widths observed experimentally were all on the order of $300 \text{ m}\text{\AA}$, which is the limit of the spectrometer resolution. No evidence for Stark broadening beyond this limit was seen, and other high density effects such as the presence of $n=2$ dielectronic satellites and continuum lowering were also absent. Code simulations showed that such effects are beyond the resolving power of the instrument for densities below about four times critical, $7.2 \times 10^{22} \text{ cm}^{-3}$, and this puts an upper limit on the density from which the observed lines were emitted. This is supported by hydrocode simulations of the experiment which will be discussed later.

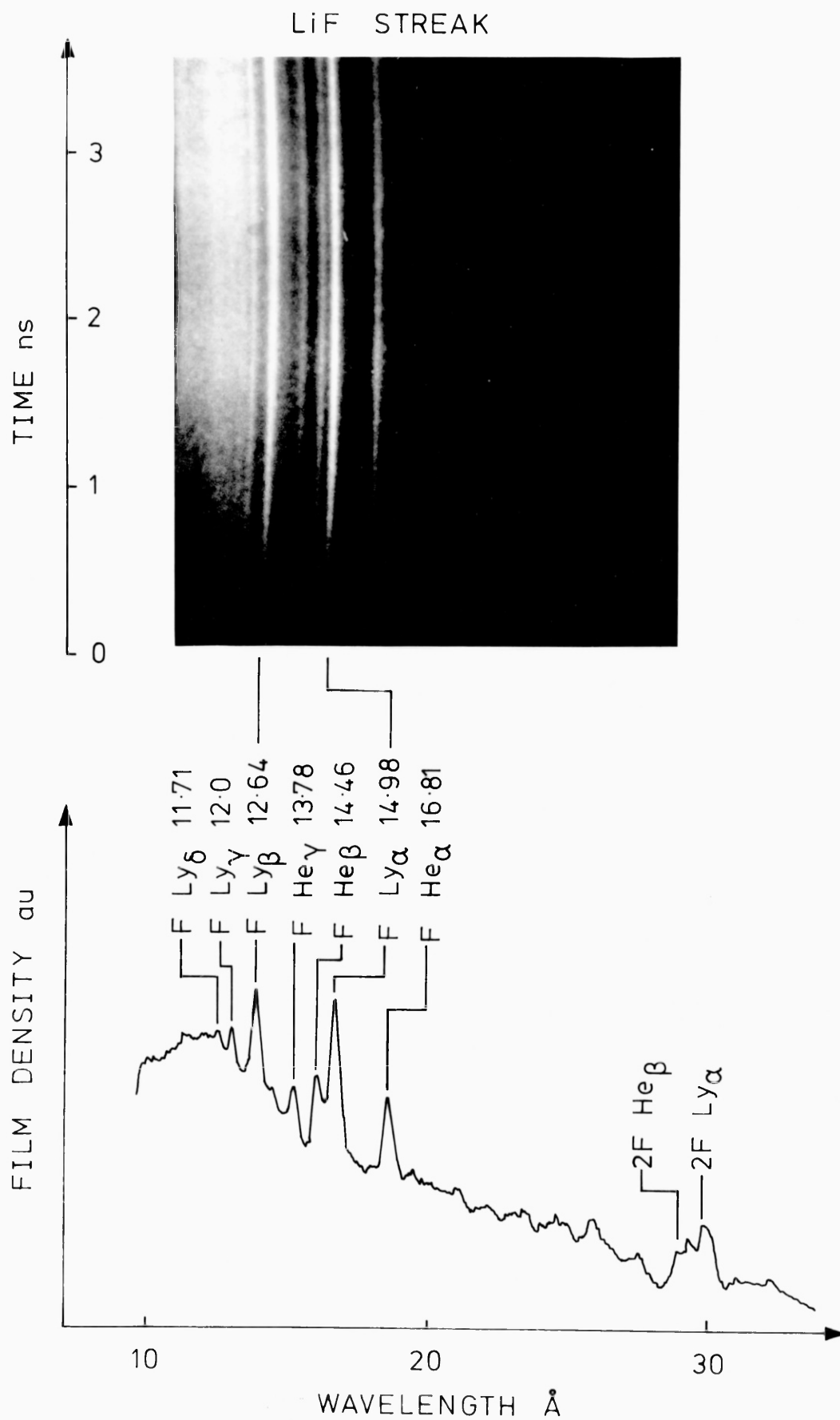


Figure 7.2 X-ray streak spectrum and densitometer trace from a LiF slab target.

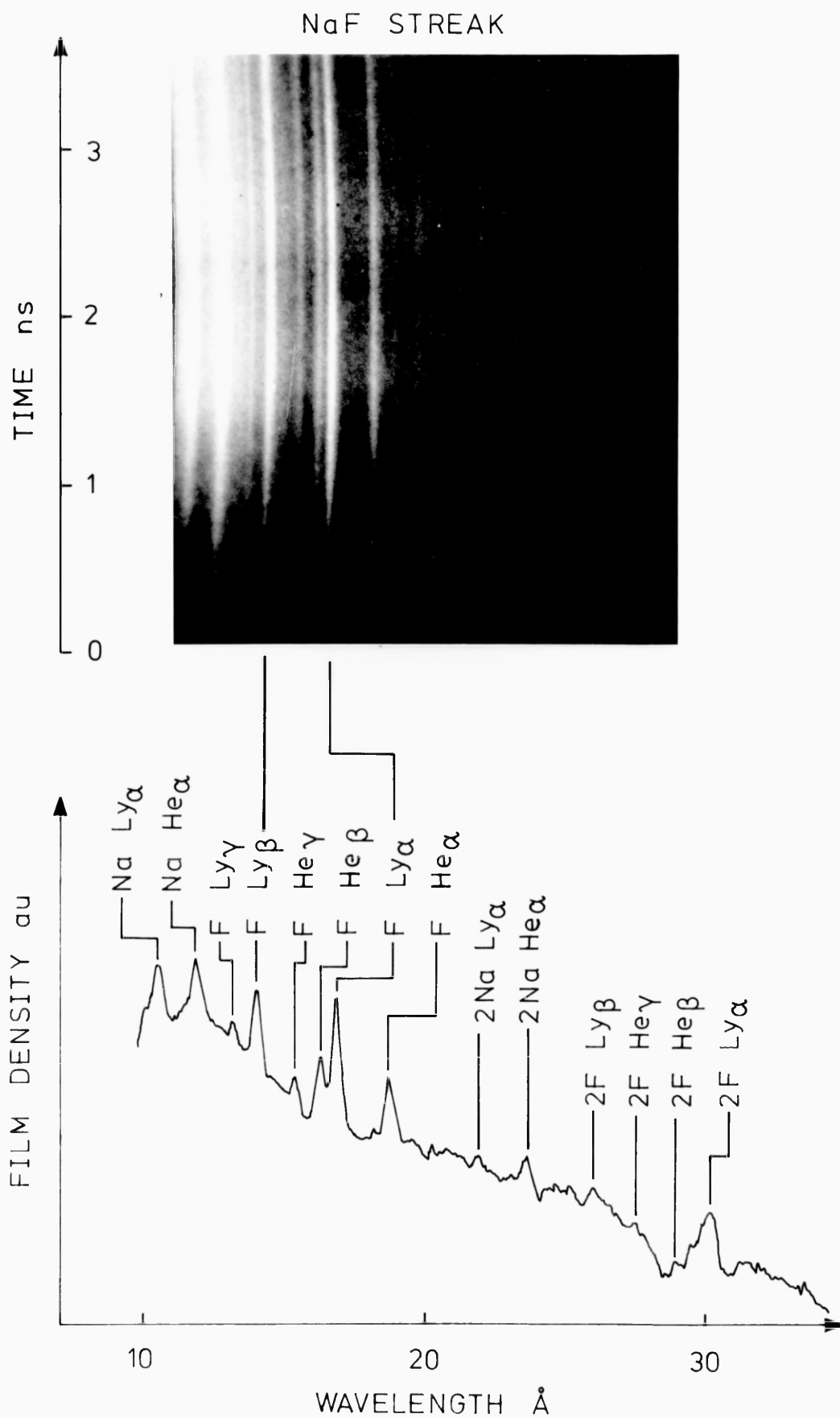


Figure 7.3 X-ray streak spectrum and densitometer trace from a NaF slab target.

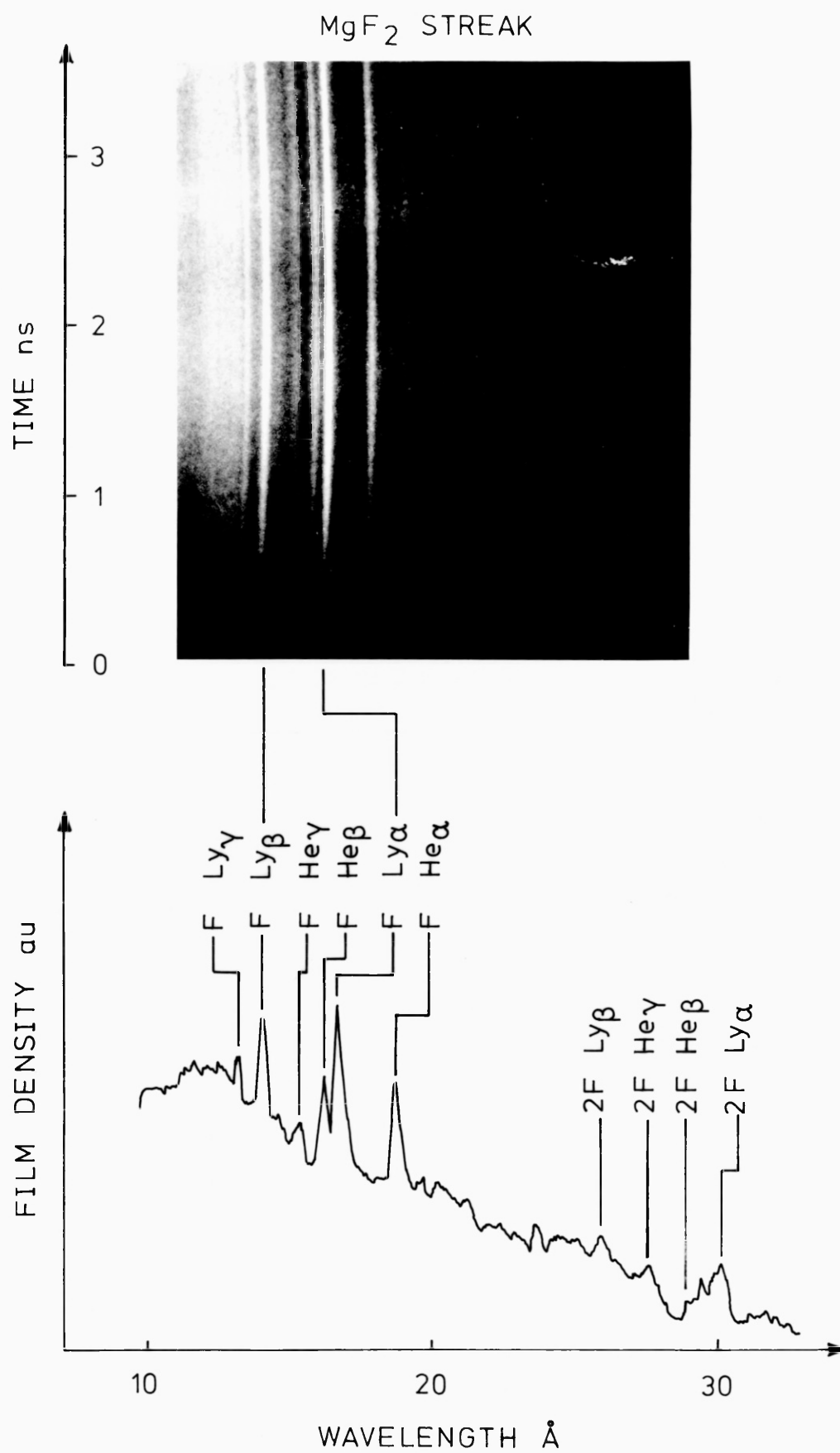


Figure 7.4 X-ray streak spectrum and densitometer trace from a MgF_2 slab target.

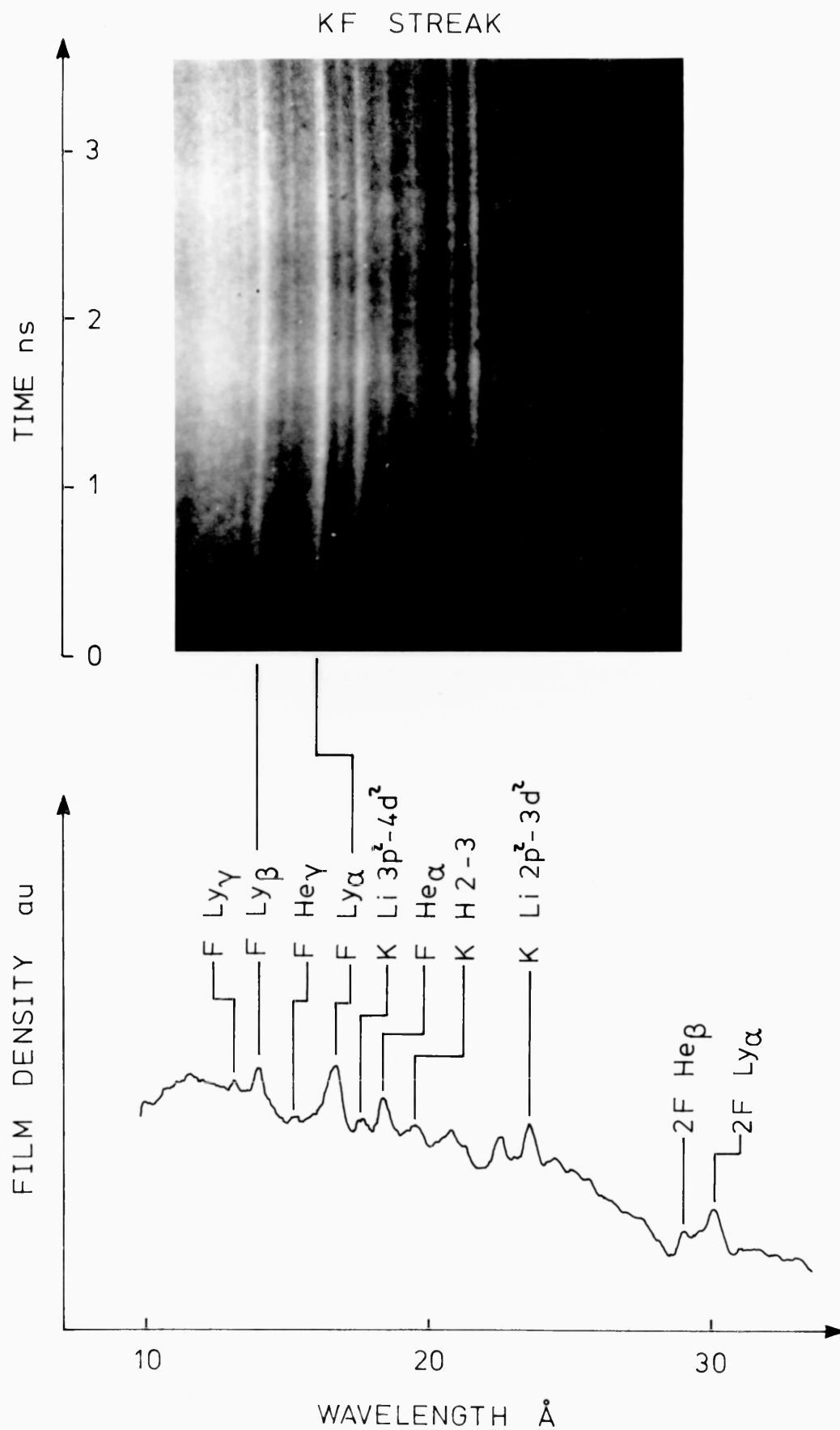


Figure 7.5 X-ray streak spectrum and densitometer trace from a KF slab target.

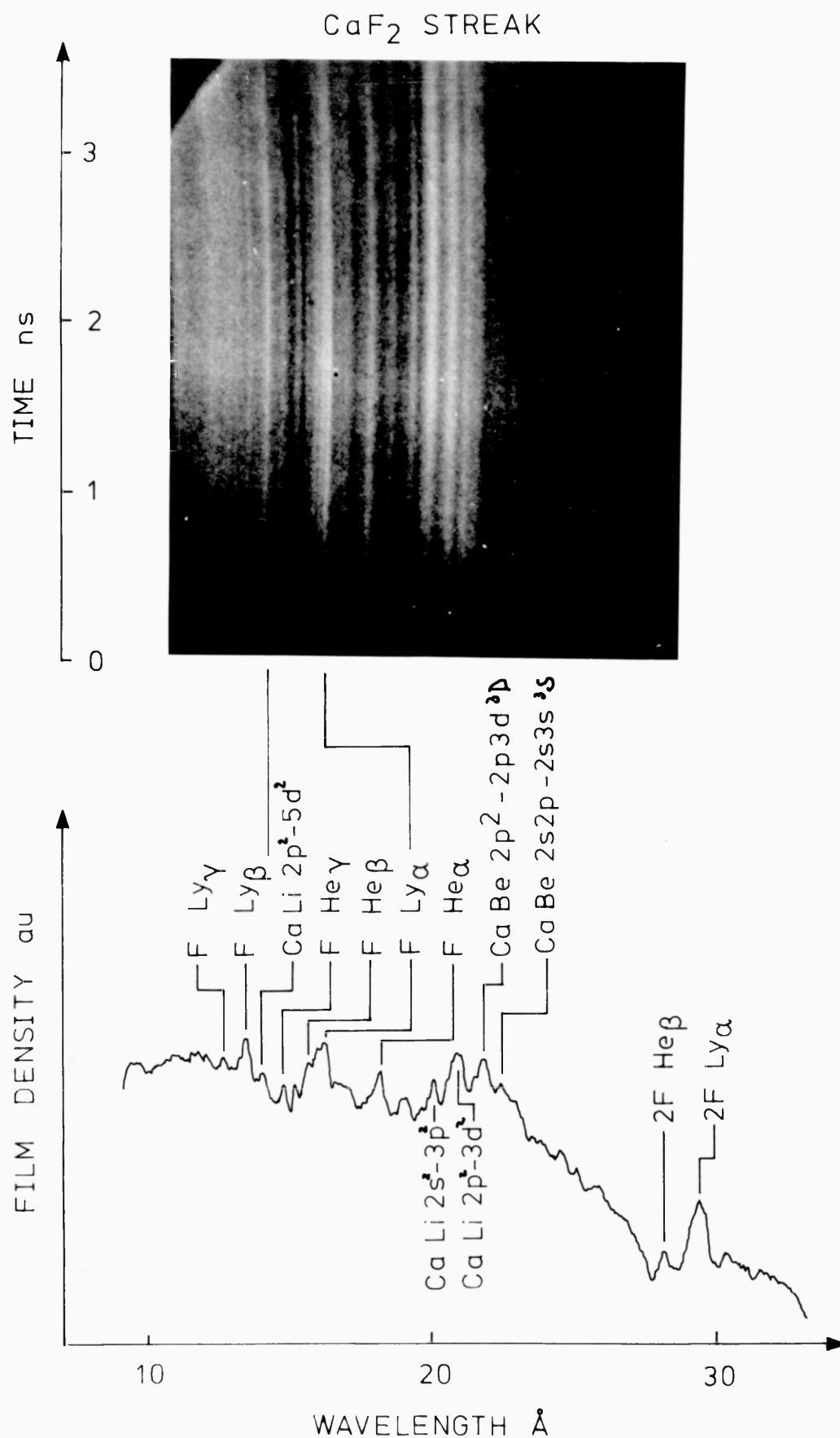


Figure 7.6 X-ray streak spectrum and densitometer trace from a CaF_2 slab target.

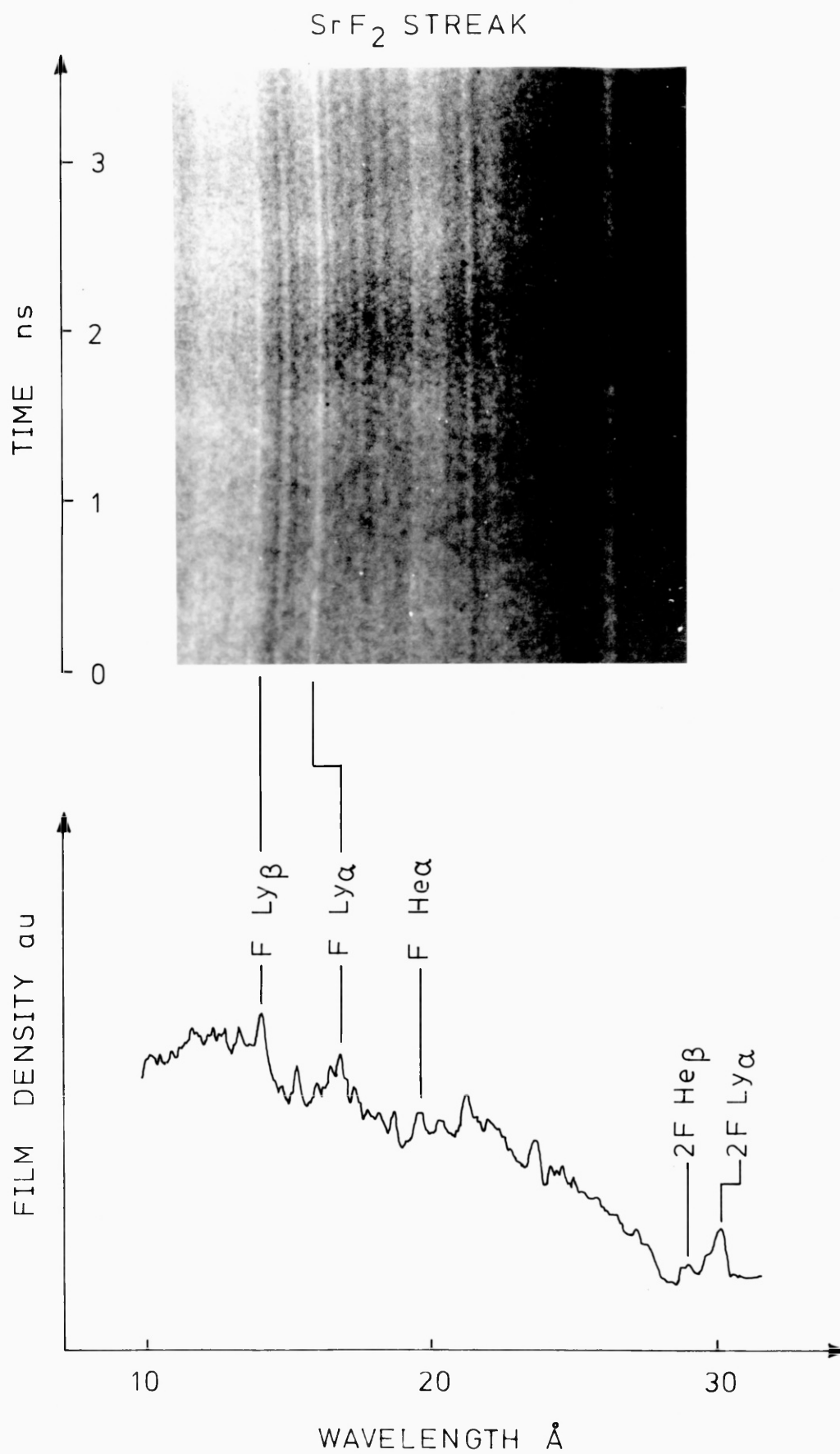


Figure 7.7 X-ray streak spectrum and densitometer trace from a SrF_2 slab target.

7.4 DETERMINATION OF THE ELECTRON TEMPERATURE

To investigate the effect of high Z dopants on the cooling rates of the plasmas produced in the experiment, the electron temperature as a function of time was found from the ratio of F Ly $_{\gamma}$ to F He $_{\gamma}$. The γ lines were chosen in preference to the stronger α and β lines as they were more likely to be optically thin, having shorter wavelengths and lower absorption oscillator strengths. This follows from an expression for the optical depth τ , which equation (2.40) gives as:—

$$\tau = \lambda^2 g(\nu) A_{n1} N_1 R$$

where λ is the wavelength of the transition, $g(\nu)$ the line profile function, A_{n1} the n to 1 transition absorption coefficient, N_1 the ground state population density and R the path length through the plasma.

Higher lying levels are also more nearly in equilibrium with the free electrons, as they are further up the recombination cascade where level populations are more likely to be dominated by collisional rather than radiative processes. However F Ly $_{\delta}$, He $_{\delta}$ and the higher series members were not well observed for all target materials at all times in this experiment and could not be used to determine the electron temperature for the full data set.

Streak spectra were traced every 1 mm (225 ps) on film along the time axis using a 1D densitometer, and the film curve deconvoluted from the data by referencing film density to a step wedge printed onto the film, to give relative intensities. The background intensity was subtracted off the Ly $_{\gamma}$ and He $_{\gamma}$ lines and the differential transmittance of the Al filter (56.5 % for Ly $_{\gamma}$ and 44.0 % for He $_{\gamma}$) taken into account. Over the short wavelength range separating these two lines the spectrometer and streak camera responses were assumed to be constant.

The Ly $_{\gamma}$ and He $_{\gamma}$ intensities were plotted as a function of time and fitted to smooth curves. Any single line intensity measurement at a particular time was subject to a large error as both Ly $_{\gamma}$ and He $_{\gamma}$ were small compared to the background on which they were superimposed. However a smooth curve could generally be found for the 14 – 16 data points from any given streak. Values of line intensity were then taken from these smooth curves to find the Ly $_{\gamma}$ to He $_{\gamma}$ line ratio as a function of time. These values were then compared to line ratios generated by the codes RATION and SPECTRA to find the temperature.

The atomic physics code RATION calculates populations of excited states for a particular ion in a collisional-radiative model, for given values of electron temperature T_e and density n_e . The code uses a steady state, non LTE model, and includes detailed bare ion, H, He and Li like states explicitly. Lower ionisation stages are also calculated and are assumed to be in LTE with the Li like ground state. The level populations produced by RATION can be processed by the code SPECTRA which calculates the H and He like resonance spectrum, including the $n=2$ dielectronic satellites. The code includes both Doppler and Stark broadening, Stark being calculated from microfields. The instrumental broadening (in this case ≈ 2 eV) can also be included. Reabsorption is calculated assuming that the plasma is a homogeneous slab of a given thickness.

RATION can be run for any element with $Z \leq 26$. However the version of SPECTRA presently implemented on the Rutherford computer system was written for the simulation of laser fusion experiments and only supports data sets for a limited number of elements. The code does not include the necessary information to generate fluorine spectra directly, but it can calculate oxygen ($Z = 8$) and neon ($Z = 10$) spectra. These two elements are very close in Z to fluorine, and an average of the values they produce is likely to be a good approximation for fluorine.

RATION and SPECTRA were run for a wide range of electron temperatures from 50 to 350 eV and electron densities from quarter critical ($4.5 \times 10^{21} \text{ cm}^{-3}$) to 4 times critical ($7.2 \times 10^{22} \text{ cm}^{-3}$). A range of plasma thicknesses from 0.1 to 500 μm were investigated, and no significant difference in line ratios was found for plasmas up to about 200 μm deep. The code was generally run for a 1 μm thick plasma, a value which was computationally convenient. For each run the ratio of Ly_γ to He_γ was found and plotted as a function of T_e . A median value for the fluorine Ly_γ to He_γ ratio was then estimated from the plots for oxygen and neon. Figure 7.8 shows the γ line ratios as a function of electron temperature T_e generated by RATION and SPECTRA for oxygen and neon, and the inferred curve for fluorine at critical electron density $n_c = 1.8 \times 10^{22} \text{ cm}^{-3}$ for 0.25 μm light. Runs at $\frac{1}{4} n_c$ and $4 \times n_c$ gave similar results.

The code SPECTRA can also generate opacities as a function of wavelength. As a check that the plasma was optically thin for the lines used in the temperature analysis, the code was run for a temperature of 200 eV and density of $1.8 \times 10^{22} \text{ cm}^{-3}$. These values were typical of those measured experimentally and suggested by hydrocode simulations.

By averaging results for neon and oxygen γ lines, values of $0.038 \mu\text{m}^{-1}$ and $0.049 \mu\text{m}^{-1}$ were found for fluorine Ly_γ and He_γ respectively. Hydrocode simulations with MEDUSA suggest that most of the X-ray emission observed comes from a few microns of plasma near critical density and propagates through at most a few hundred microns of coronal plasma which is fully ionised and underdense. In this case the fluorine lines of interest will have optical depths well below 1, and will be optically thin.

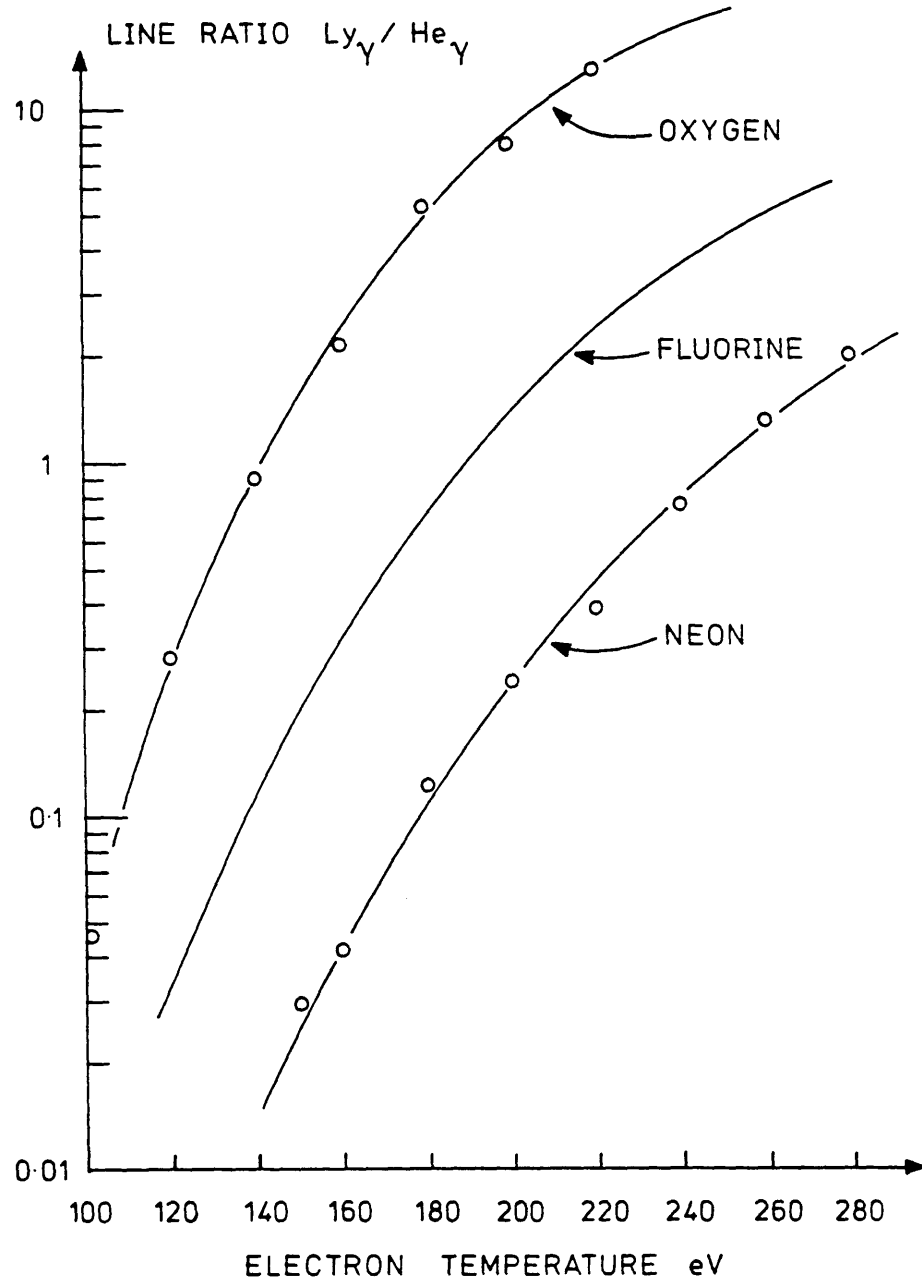


Figure 7.8 Ly_γ to He_γ line ratios as a function of electron temperature T_e generated by RATION and SPECTRA for oxygen and neon, and the inferred curve for fluorine.

7.5 EXPERIMENTAL COOLING RATES

Figures 7.9 and 7.10 show the measured time histories of $F Ly_{\gamma}$ and He_{γ} emission from a LiF slab target irradiated at an intensity of $1.5 \times 10^{14} \text{ Wcm}^{-2}$. At early times it is relatively easy to fit a smooth curve to the line intensity, but in both cases the data becomes rather noisy at later times as the line intensity becomes smaller relative to the background on the streak.

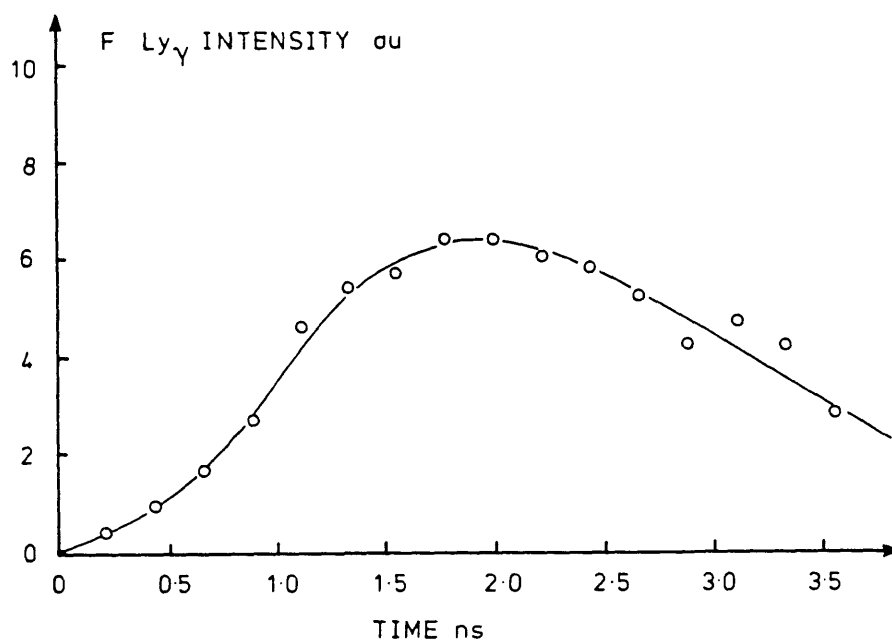


Figure 7.9 Experimental time evolution of $F Ly_{\gamma}$ emission from a LiF target.

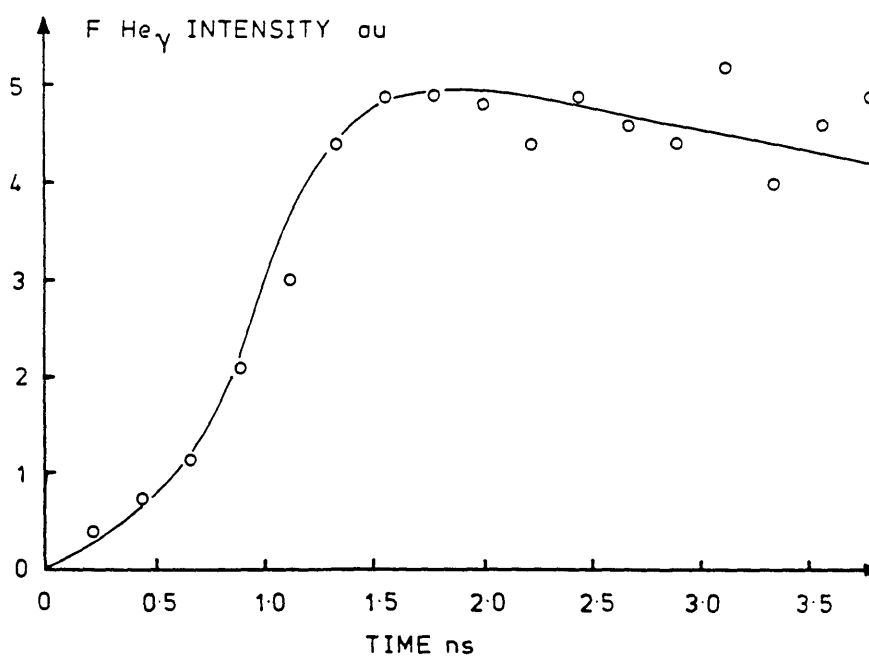
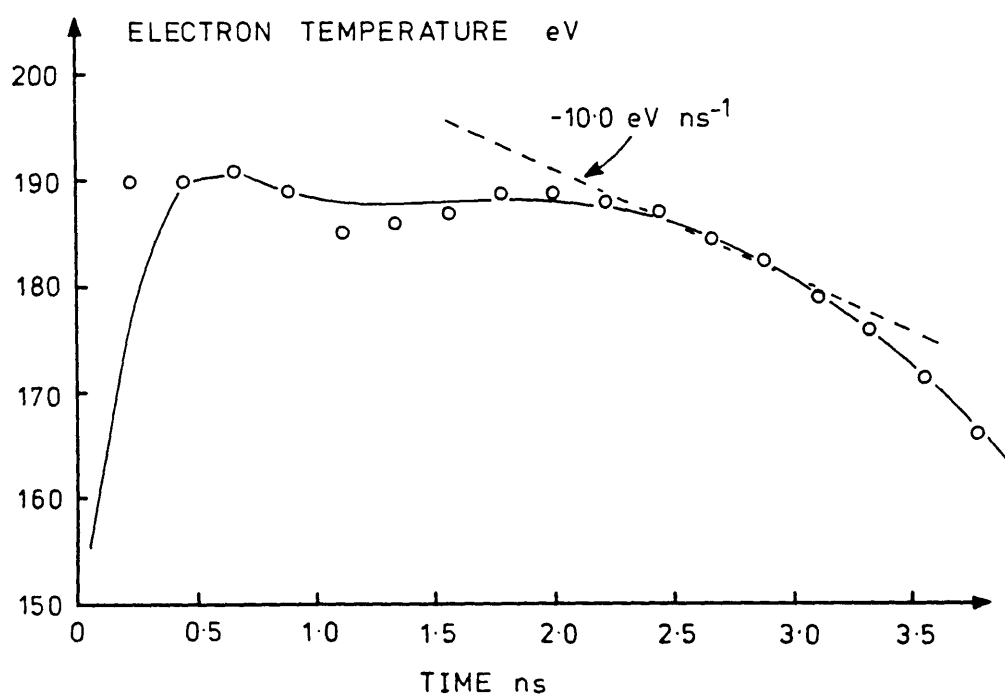


Figure 7.10 Experimental time evolution of $F He_{\gamma}$ emission from a LiF target.

As described in section 7.4 the electron temperature as a function of time was found by taking line intensities from the smooth curves and comparing experimental $F Ly_{\gamma}$ to He_{γ} ratios to those predicted by figure 7.8. Figures 7.11 to 7.15 show the experimental time evolution of electron temperature for LiF, NaF, MgF_2 , KF and CaF_2 targets found using this technique. Due to the large number of line overlaps, strong background and small signal size in the case of the SrF_2 spectrum, it was however impossible to accurately follow the temperature for this target.



7.11 Experimental time evolution of T_e for a LiF slab target.

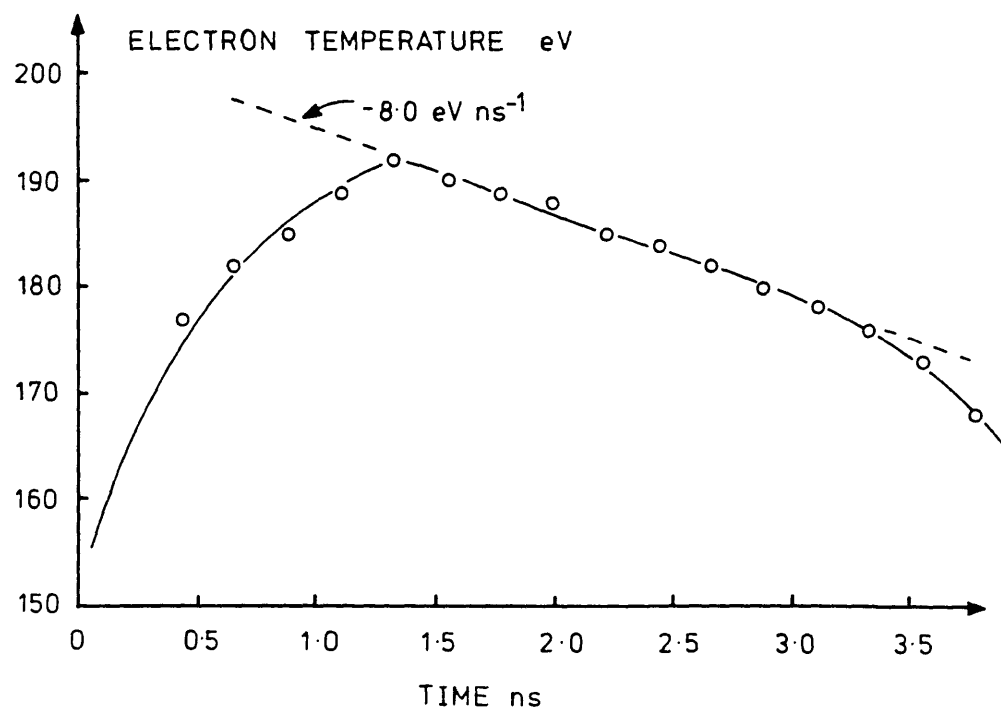


Figure 7.12 Experimental time evolution of T_e for a NaF slab target.

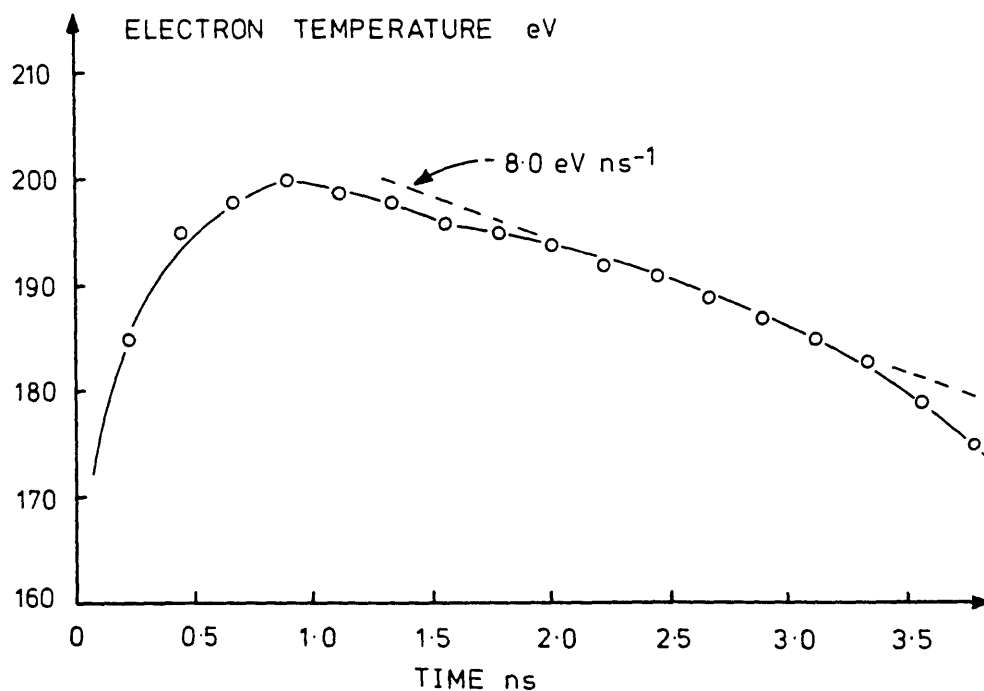


Figure 7.13 Experimental time evolution of T_e for a MgF_2 slab target.

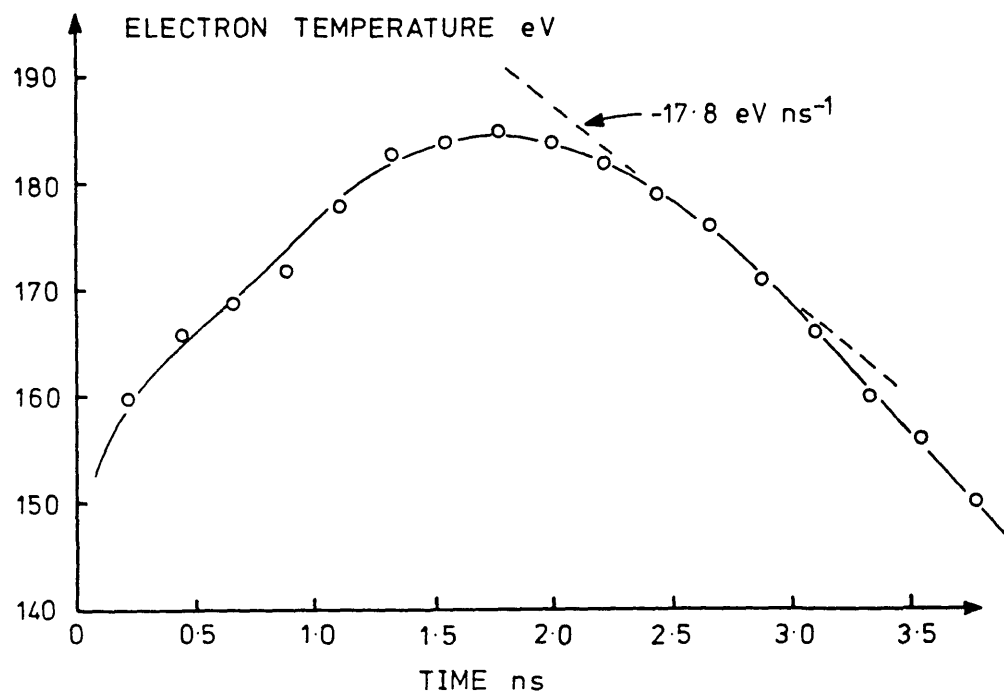


Figure 7.14 Experimental time evolution of T_e for a KF slab target.

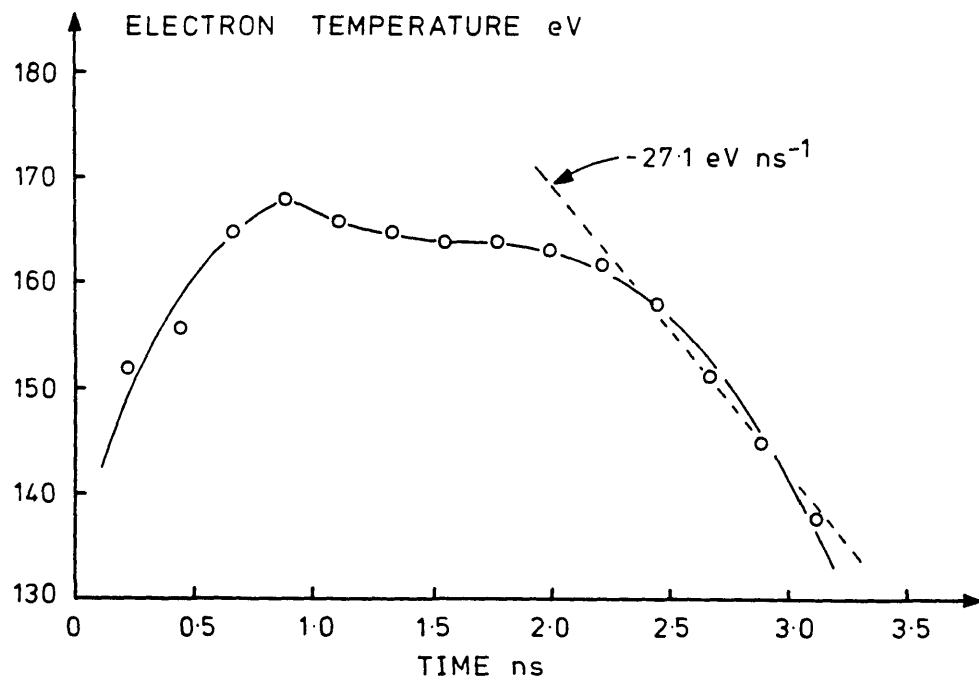


Figure 7.15 Experimental time evolution of T_e for a CaF_2 slab target.

The time histories of electron temperature for LiF, NaF and MgF₂ show quite similar profiles, with peak temperatures in the 190–200 eV range near the peak of the laser pulse at 1 ns. KF and CaF₂ targets however show a less rapid rise in temperature during the rising edge of the laser pulse, reflecting the additional energy required to heat and ionise the high Z dopants. In addition there is a more rapid cooling of the plasma after the turn off of the laser at around 2 ns.

The cooling rate for these targets was quantified by carrying out a linear least squares fitting to the electron temperature from 0.25 to 1.15 ns after the end of the laser pulse at 2 ns. This corresponds to the sort of timescale over which population inversions and gain were seen in fibre based X-ray laser experiments. LiF, NaF and MgF₂ targets were found to have similar cooling rates of -10.0 , -8.0 and -8.0 eV ns⁻¹ respectively. This is not a surprising result, considering the similar atomic numbers of fluorine ($Z = 9$), sodium ($Z = 10$) and magnesium ($Z = 11$).

In contrast the KF and CaF₂ targets showed cooling rates of -17.8 and -27.1 eV ns⁻¹ over the same time interval, an enhancement by a factor of 2 or 3 over the cooling rates for the lower Z targets. These results are summarised in figure 7.16 which is a plot of measured cooling rate against the highest atomic number in the target.

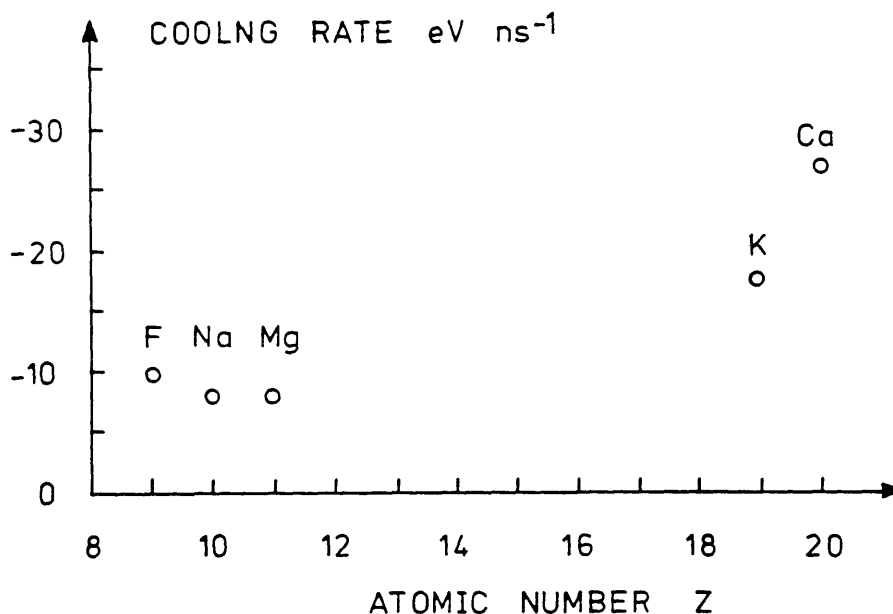


Figure 7.16 Plasma cooling rate eV ns⁻¹ versus atomic number of dopant.

7.6 HYDRODYNAMIC SIMULATIONS.

The 1D Lagrangian hydrocode MEDUSA was used to simulate the heating phase of the experiment during the laser pulse. This was done as a check that experimental values of electron temperature were sensible, and to provide some idea of the plasma density profile for inclusion in the atomic physics modelling.

The code was used to model a $0.25 \mu\text{m}$, 2 ns FWHM Gaussian laser pulse with a peak power of $1.5 \times 10^{14} \text{ Wcm}^{-2}$ normally incident on a solid fluorine slab target of density 2.64 g cm^{-3} . This should give a reasonable approximation to a LiF slab target, as lithium is a low $Z = 3$ element and should contribute little to the plasma mass or energy required to heat the target.

Laser energy was assumed to be absorbed by inverse Bremsstrahlung in the coronal plasma up to critical density at $n_c = 1.8 \times 10^{22} \text{ cm}^{-3}$. At the critical density surface a fraction 0.05 of the remaining laser energy (that not absorbed by inverse Bremsstrahlung) was dumped to simulate anomalous absorption. 90 % of this energy was placed in "hot" electrons with a semi-empirical temperature distribution (a function of $I\lambda^2$) due to Giovanelli et-al [92]. The thermal conductivity of the electrons was set to the Spitzer value, but a flux limiter of 0.033 times the classical free streaming limit was imposed to limit the heat flow in steep temperature gradients. This is an empirical value which has been found to give good fits for a number of experiments of this type [93].

Figure 7.17 shows the output of MEDUSA at the peak of the laser pulse. Electron temperature T_e , electron density n_e and the average ionisation stage Z are plotted as a function of distance from the initial target surface. At this point in time the heat front has propagated some $1.5 \mu\text{m}$ into the target. The plasma has expanded out to some $150 \mu\text{m}$, but beyond $7 \mu\text{m}$ the plasma is over heated and fully stripped. The density also falls quite rapidly in this region and it is likely that the large coronal plasma will contribute little to the emission or opacity.

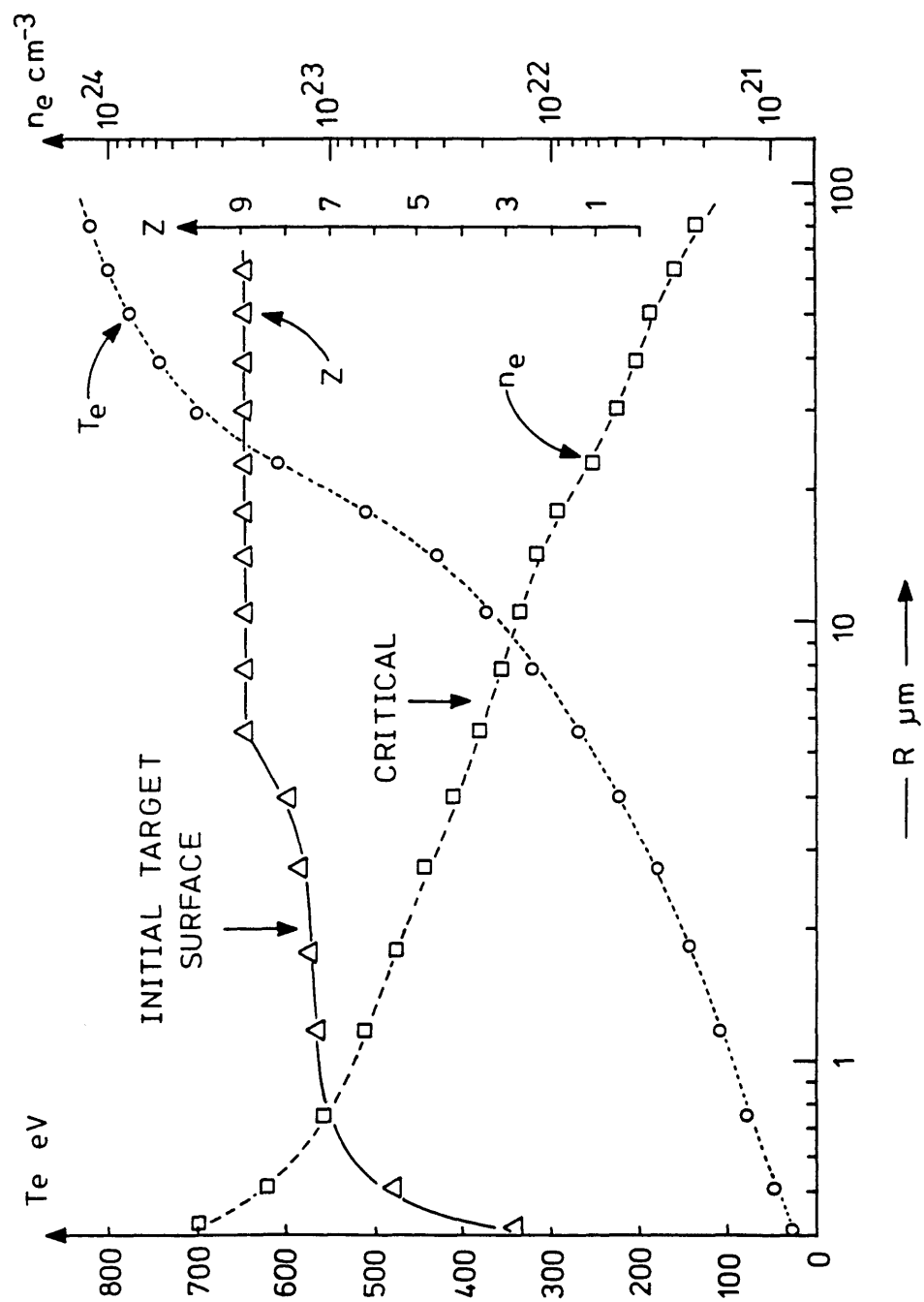


Figure 7.17 MEDUSA hydrocode simulation of a solid fluorine slab target irradiated by a 2 ns pulse of $0.25 \mu\text{m}$ light, with a peak power of $1.5 \times 10^{14} \text{ Wcm}^{-2}$.

7.7 CONCLUSION

A range of solid fluoride slab targets were irradiated in point focus geometry with the SPRITE laser system to investigate the effect of high Z dopants on recombining H like fluorine plasmas. Time resolved X-ray spectroscopy was used to diagnose the plasmas and the time evolution of electron temperature was found by comparing measured F Ly $_{\gamma}$ to He $_{\gamma}$ ratios with predictions by the atomic physics codes RATION and SPECTRA. Temperatures measured experimentally using this technique were found to be in good agreement with MEDUSA hydrocode simulations during the heating phase. The cooling phase of the experiment could not be simulated with the present version of the code as it does not include radiation transport. A version which does include radiation is presently being used by the author and others to model the cooling of these plasmas.

After the turn off of the laser pulse, LiF, NaF and MgF $_2$ targets showed comparable cooling rates of -10.0 , -8.0 and -8.0 eV ns $^{-1}$, as might be expected from the similar atomic numbers of fluorine, sodium and magnesium. In contrast targets containing potassium and calcium were found to have cooling rates of -17.8 and -27.1 eV ns $^{-1}$, a factor of some 2 or 3 times higher than the low Z targets. This enhanced cooling is believed to be due to increased radiation loss from plasmas containing high Z dopants, as radiation loss by Bremsstrahlung and radiative recombination scales as Z^2 while line radiation loss scales as Z^4 . Plasmas containing strontium however were found to be over cooled and contained many line overlaps with the fluorine Ly and He series.

The enhanced cooling of hydrogenic fluorine plasmas by high Z dopants observed experimentally might well result in the formation of population inversions at higher densities in X-ray laser experiments. Radiative cooling might also be applied to shorter wavelength recombination X-ray laser schemes in order to reduce the strong scaling of initial plasma density with wavelength needed to give significant gain.

The pulse length of 2 ns used in this experiment was rather long compared to the timescale over which we wished to investigate the cooling of laser produced plasmas. It is likely that the cooling effect of high Z dopants would be more pronounced and more easily observed with laser pulse lengths in the tens of picoseconds regime, however such a pulse length was unavailable at the time of this investigation. It is hoped that a more complete investigation of this effect with a shorter pulse length will be carried out in the future.

CHAPTER 8. PICOSECOND SPECTROSCOPY

ABSTRACT

The strong Z dependence of initial density in the fibre based recombination amplifier places a serious limitation on the short wavelength scalability of this scheme. However short pulse irradiation of solid targets with picosecond pulses can provide the correct starting conditions for a hydrogen like aluminium recombination water window amplifier.

The interaction of single 3.5 ps, 2.5 J pulses of KrF light with solid targets at irradiances up to 10^{17} W cm $^{-2}$ was studied with time resolved and time integrated X-ray and XUV spectroscopy. Fully stripped Al plasmas at temperatures of ≈ 400 eV were observed, and densities close to solid inferred from X-ray line Stark widths and continuum lowering. These results were well modelled by 1D hydrodynamic and 2D Fokker-Planck codes, which also showed that lateral heat transport is not important on the picosecond timescale.

8.1 INTRODUCTION

The interaction of short (picosecond and below) high intensity laser pulses with matter is highly topical. In the picosecond regime the laser plasma interaction takes place on a timescale short compared to that on which significant hydrodynamic motion can take place. This can result in the formation of hot, high density plasmas which are of great interest for basic plasma and atomic physics studies, as well as for the scaling of recombination X-ray laser schemes to shorter wavelengths. Until recently plasmas at near solid densities have only been produced in the laboratory in spherical compression experiments. These are both difficult to diagnose and expensive in terms of the number of compression beams and associated optics required.

For long (> 20 ps) pulses the leading edge of the pulse ablates a plasma which has time to expand away from the target during the interaction. The majority of the beam interacts with the preformed plasma via inverse Bremsstrahlung at densities below critical, and by resonance processes at the critical surface. Energy is then transported by various mechanisms through the plasma to the target.

In the picosecond regime the laser plasma interaction is qualitatively quite different. The plasma sound speed [94] is given by $c_s \approx (\zeta Z_i K T_i / m_i)^{1/2}$, where ζ is the adiabatic index, Z_i the ion charge, K Boltzmann's constant, T_i the ion temperature and m_i the ion mass. For say a 400 eV fully stripped Al plasma c_s will be on the order of 10^7 cm s^{-1} . If the hydrodynamic expansion velocity of the plasma is on the same order as the sound speed then clearly the plasma will expand less than a micron away from the target surface during a picosecond laser pulse. With very little underdense plasma present, inverse Bremsstrahlung absorption becomes inefficient. This leads to a large fraction of the coupled laser energy being dumped close to the target surface at critical density and transported rapidly to the solid region of the target. The absorption mechanism itself is not well understood and various types of non-collisional processes such as "non classical" resonance absorption [95], resonance absorption at $4 \times$ critical in very steep density gradients [96,111] and anomalous skin-effects [97] have been suggested. However the pulse length in this experiment is believed to be long enough to generate sufficient overdense plasma for "classical" resonance absorption to take place [98].

The "classical" resonance absorption process requires oblique incidence light and the first harmonic drives longitudinal plasma waves which accelerate electrons. Whatever the absorption mechanism, hot, high density plasmas can be generated by picosecond pulses in planar geometry experiments which are much easier to perform and diagnose.

Other workers in the picosecond regime [99,100,101,110] have used relatively low energy (20 mJ) pulses and integrated over several thousand shots to record one result. Their laser systems also generated significant ASE prepulses which broke down the target surface and resulted in the short pulse interacting with a large underdense preformed plasma at some distance from the initial target surface. This chapter reports on the first single shot picosecond KrF interaction study, and was carried out with prepulse levels lower than the breakdown threshold of the target to give a direct laser / target interaction.

8.2 THE SPRITE SHORT PULSE KrF SYSTEM

The SPRITE laser at RAL is a modular, developmental high power KrF system which can be reconfigured for different applications. This experiment was carried out with a recently developed short pulse set up [102], see figure 8.1. 1 nJ, 2.5 ps seed pulses of 0.746 μm light were produced in a commercial optical dye oscillator pumped by a frequency doubled, actively mode locked CW Nd:YAG laser. Single pulses were switched out of the oscillator, amplified in a 3 stage XeCl excimer pumped dye amplifier chain, and frequency tripled to 0.248 μm in two KDP crystals. Frequency tripled quarter micron 1 μJ pulses were then spatially filtered and used to seed a chain of 3 KrF amplifiers. The first pre-amplifier was a commercial 8 Hz discharge system with a single pass gain of ≈ 100 . The majority of the ASE background on which the final short pulse output was superimposed originated in this amplifier. The second pre-amplifier GOBLIN was a large aperture (60 mm) e-beam pumped system built in house, and was double passed to give a gain of ≈ 200 . The final amplifier was the 270 mm aperture e-beam pumped SPRITE module with a double pass gain of ≈ 130 . By careful optimisation of the gain staging and gas fill in the amplifiers the prepulse level could be minimised, and a 3.5 ps, 2.5 J short pulse output superimposed on a 20 ns ASE background of less than 25 mJ was regularly achieved [102]. This gave corresponding short pulse powers of up to 8×10^{11} W and background powers of $\leq 1.3 \times 10^6$ W.

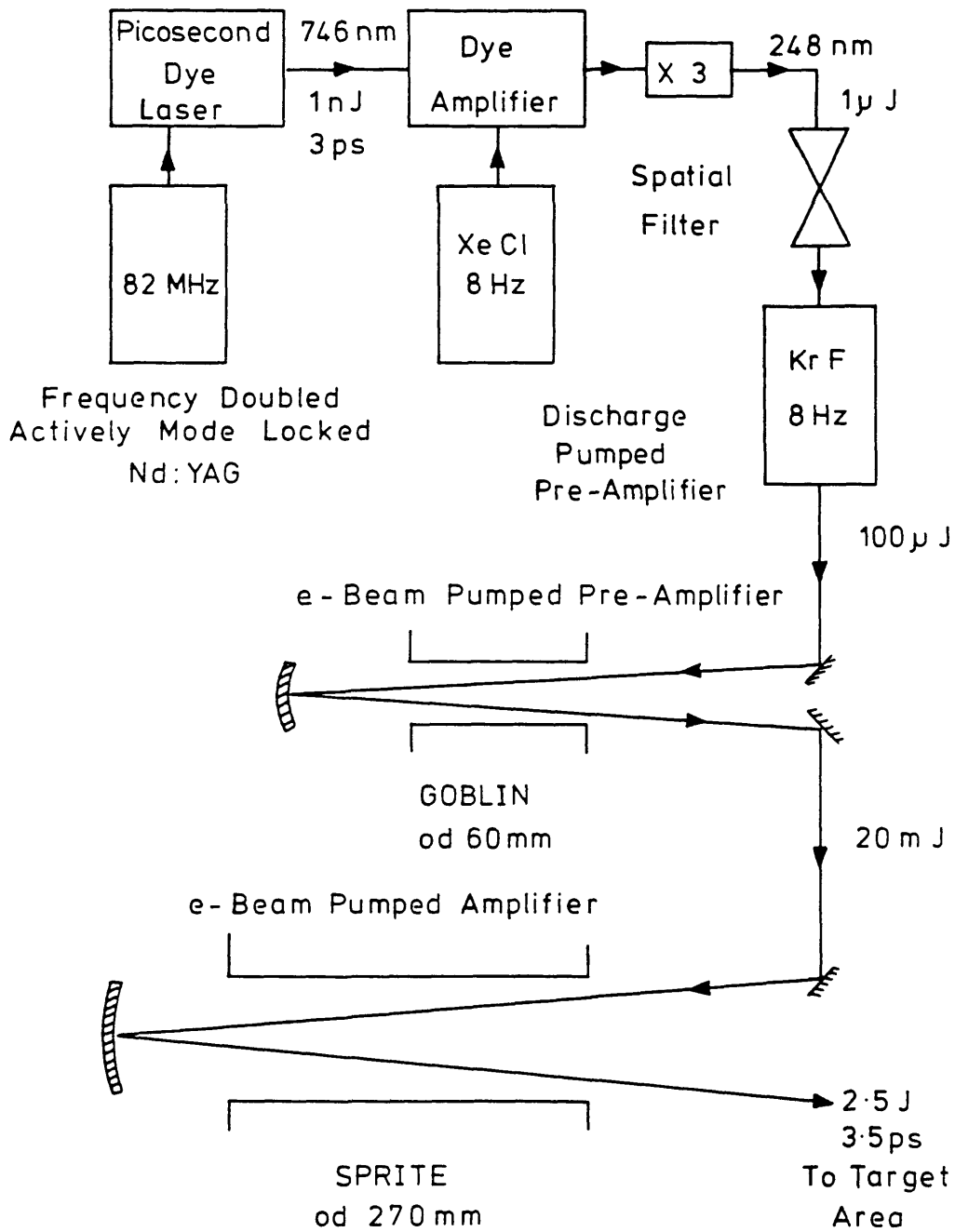


Figure 8.1 The SPRITE short pulse, high power KrF laser system.

This high contrast between prepulse and short pulse powers resulted in prepulse irradiances of $< 10^5$ of the short pulse on target, even before taking into account the poor focusing of the ASE signal. As the ASE originates some way along the amplifier chain it has a large divergence, and damage marks from shots specifically fired with a large ASE signal were some 200 – 300 μm in diameter (see figure 8.2), compared to the 20 μm focus of the short pulse. Experimentally it was found that the prepulse did not break down Al targets before the arrival of the short pulse.

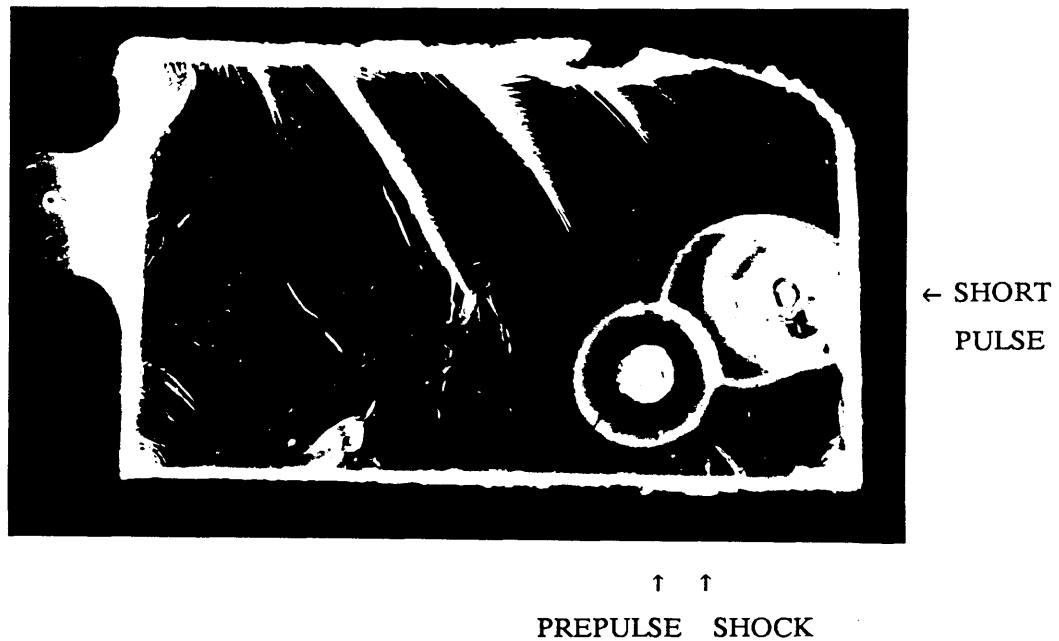


Figure 8.2 Damage marks from SPRITE pulses on a CH slab target. The two 20–30 μm spots centre right are from short pulses. The 200–300 μm spot lower right is from a shot with large prepulse. The larger circles are from shock damage.

8.3 EXPERIMENTAL

Figure 8.3 shows the experimental arrangement. Foil and slab targets were irradiated with 2.5 J, 3 ps pulses of 0.25 μm light from the SPRITE laser. An $f / 2.5$ aspheric doublet lens was used to focus the 8 cm beam to a 20 μm focal spot to give irradiances up to 10^{17} Wcm^{-2} on target. The focal spot size was monitored using active and passive pinhole cameras with 15 μm resolution while the focal spot distribution was recorded with an equivalent plane system using a similar lens. Fast optical diodes were used to measure the prepulse level and short pulse energy on each shot. Figure 8.4 shows a typical optical and X-ray diode summed trace for a 0.91 J short pulse with low ASE background.

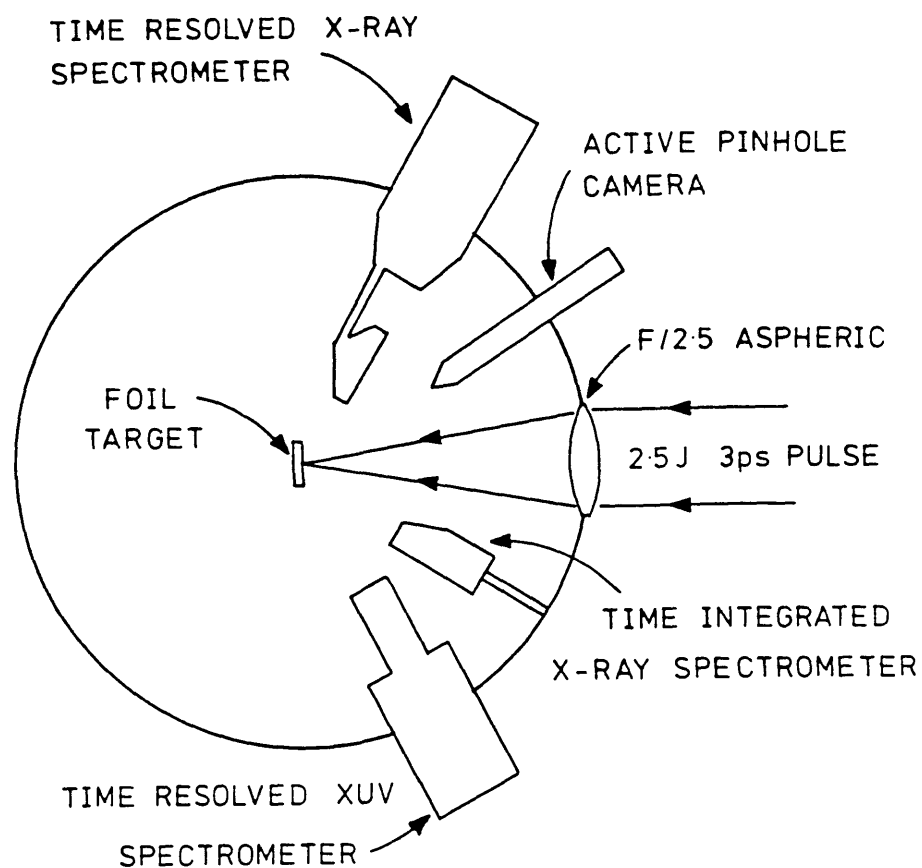


Figure 8.3 Experimental arrangement for the short pulse interaction study.

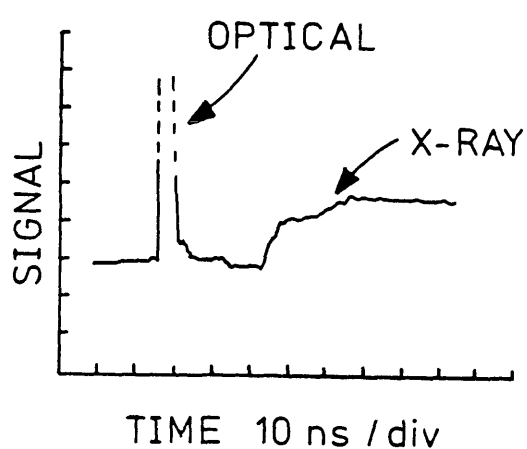


Figure 8.4 Typical optical diode and X-ray diode summed trace for a 0.91 J, 3ps short pulse, superimposed on a low ASE background.

The primary diagnostic was a time resolved X-ray spectrometer. A flat TiAP crystal ($2d = 25.76 \text{ \AA}$) was coupled to an X-ray streak camera (CsI on beryllium photocathode) which viewed the plasma at 45° to the incident laser beam. X-ray emission over a $5 - 7 \text{ \AA}$ spectral window was recorded with $15 \text{ m\AA}$ and 15 ps resolution over a 400 ps time interval. Time integrated spectra were also recorded for some shots over a similar spectral range with flat crystal, time integrated PET and Si 220 spectrometers using DEF X-ray film.

Time resolved XUV spectra were recorded with a flat field $1200 \text{ lines per mm}$ spectrometer coupled to a low magnification streak camera. This instrument had $90 - 210 \text{ \AA}$ spectral and 400 ps temporal windows, with $0.5 \text{ \AA}$ and 15 ps resolution. This spectrometer was used to observe emission due to the short pulse interaction and to look for an XUV precursor from any large ASE prepulse. The dynamic range of this instrument was sufficient to detect precursor levels some 20 times less intense than emission from the short pulse interaction. The instrument was protected from reflected laser light (to which the photocathode is very sensitive) by a $1000 \text{ \AA}$ thick Al foil filter.

8.4 RESULTS

8.4.1 X-RAY SPECTRAL OBSERVATIONS FROM LOW PREPULSE SHOTS

A number of features typical of ultra high density plasmas such as Stark broadening and continuum lowering were seen in X-ray spectra from shots with a low ASE prepulse. Figure 8.5 shows an X-ray streak spectrum from an Al foil target irradiated with a $3 \times 10^{16} \text{ Wcm}^{-2}$ short pulse and an ASE prepulse irradiance of $< 3 \times 10^{11} \text{ Wcm}^{-2}$. The Al He_β ($1s^2-1s3p$) transition is very wide ($\approx 80 \text{ m\AA}$ FWHM) and intense, completely dominating the spectrum. Higher series transitions from H and He like ions (He_γ $1s^2-1s4p$ etc) are faint and start to turn on some 15 ps after the He_β transition. The whole spectrum is short lived ($\approx 150 \text{ ps}$) and this is indicative of a small, rapidly expanding and cooling plasma.

The X-ray spectrometer has a dispersion of $\approx 100 \text{ m\AA mm}^{-1}$, and a magnification of $\approx 1:1$ at the streak camera photocathode. A $20 \text{ }\mu\text{m}$ source would give a corresponding source broadening of $\approx 2 \text{ m\AA}$, much less than the observed line width. The resolution of this instrument due to the TiAP crystal rocking curve is $\approx 12 \text{ m\AA}$, again significantly smaller than the observed line width.

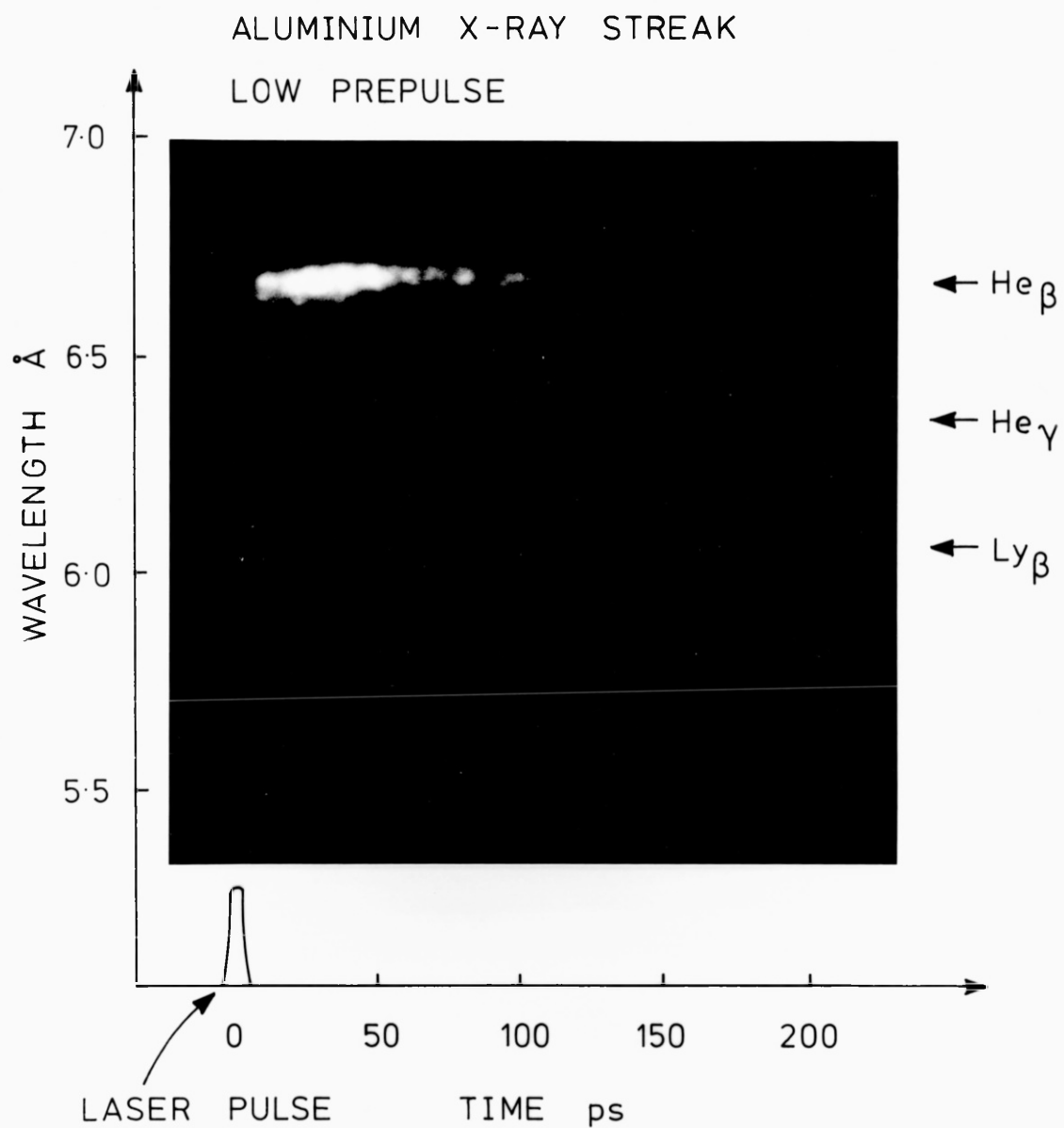


Figure 8.5 X-ray streak spectrum from an Al slab target irradiated at an intensity of $3 \times 10^{16} \text{ Wcm}^{-2}$ with a 3 ps SPRITE pulse.

Similarly the Doppler half width for a Gaussian line profile given by equation (2.20) is :-

$$\Delta\lambda_{\frac{1}{2}} \approx \lambda \left[\frac{2 K T \text{Ln}2}{m c^2} \right]^{\frac{1}{2}}$$

For a 400 eV aluminium plasma the Doppler half width for He_{β} is ≈ 1.2 mÅ (FWHM), again well below the observed width. The broadening in the X-ray spectrum is actually due to the Stark effect [103], where the microfield from a high density plasma is large enough to significantly perturb the energy levels of an emitting ion. Both the Stark width and profile of these X-ray lines are powerful diagnostics of the electron density n_e [104,105].

The suppression of high series members at early times is a related effect, continuum lowering [106,107]. The Stark widths of these lines becomes comparable to half the inter series separation and the wavefunctions for two adjacent transitions mix. These transitions become poorly resolved and merge into the continuum. At later times the density falls, lessening the Stark broadening and allowing these transitions to turn on. Continuum lowering can also be used to diagnose the density, but is much less accurate than using Stark profiles due to the difficulty in deciding when a line actually merges with the continuum.

Figure 8.6 compares densitometer traces of the spectrum in figure 8.5 taken at 20 ps and 35 ps after the laser pulse. It can clearly be seen that the Stark width of He_{β} decreases in time as the plasma expands, and that Ly_{β} turns on later than He_{β} .

Similar results were seen in time integrated X-ray spectra. Figure 8.7 shows two spectra from Al foil shots with low and high levels of ASE prepulse. In the low prepulse case no lines at wavelengths shorter than Al He_{β} are seen. Even though this spectrum is time integrated most of the emission takes place at early times when the density is high, and continuum lowering is still in effect. Ly_{α} is also very bright in this spectrum indicating that the plasma is fully ionised.

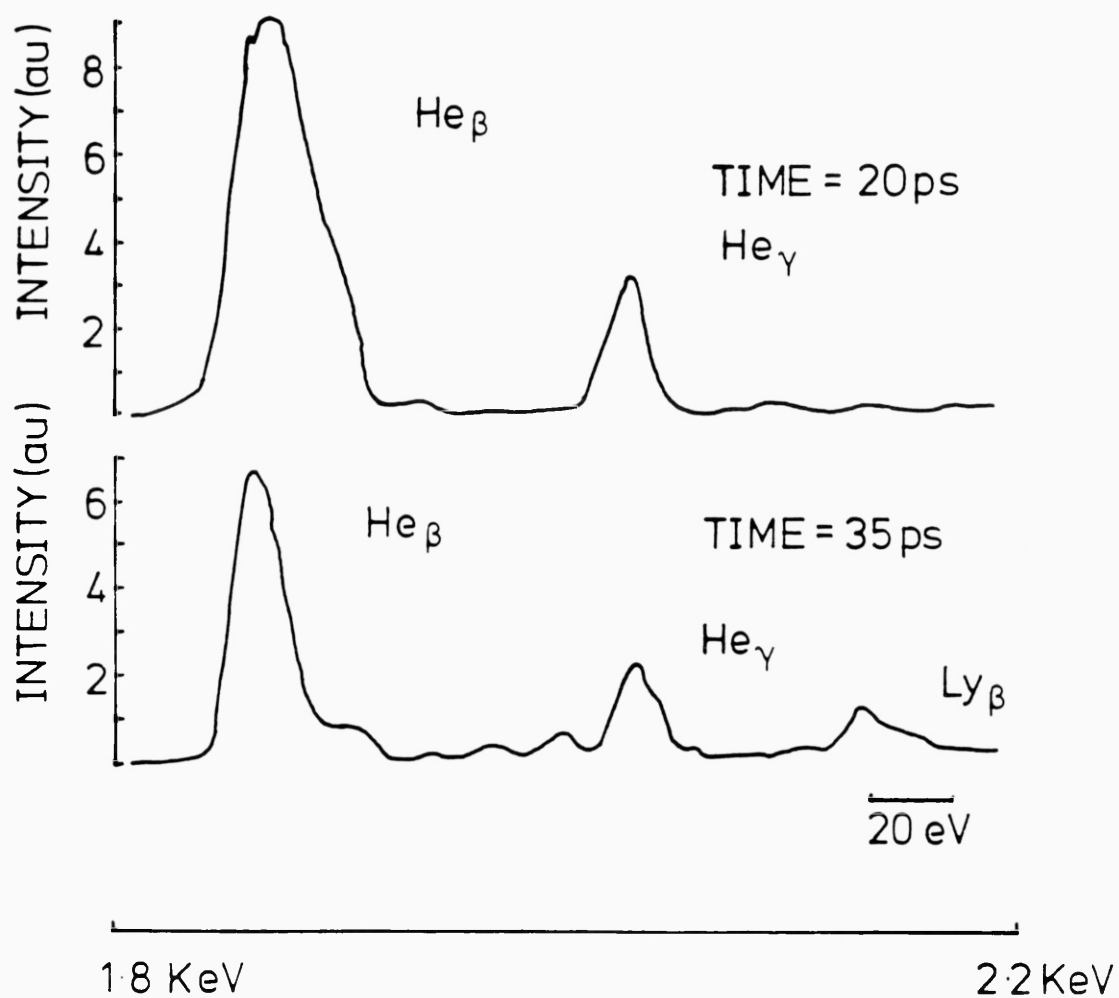
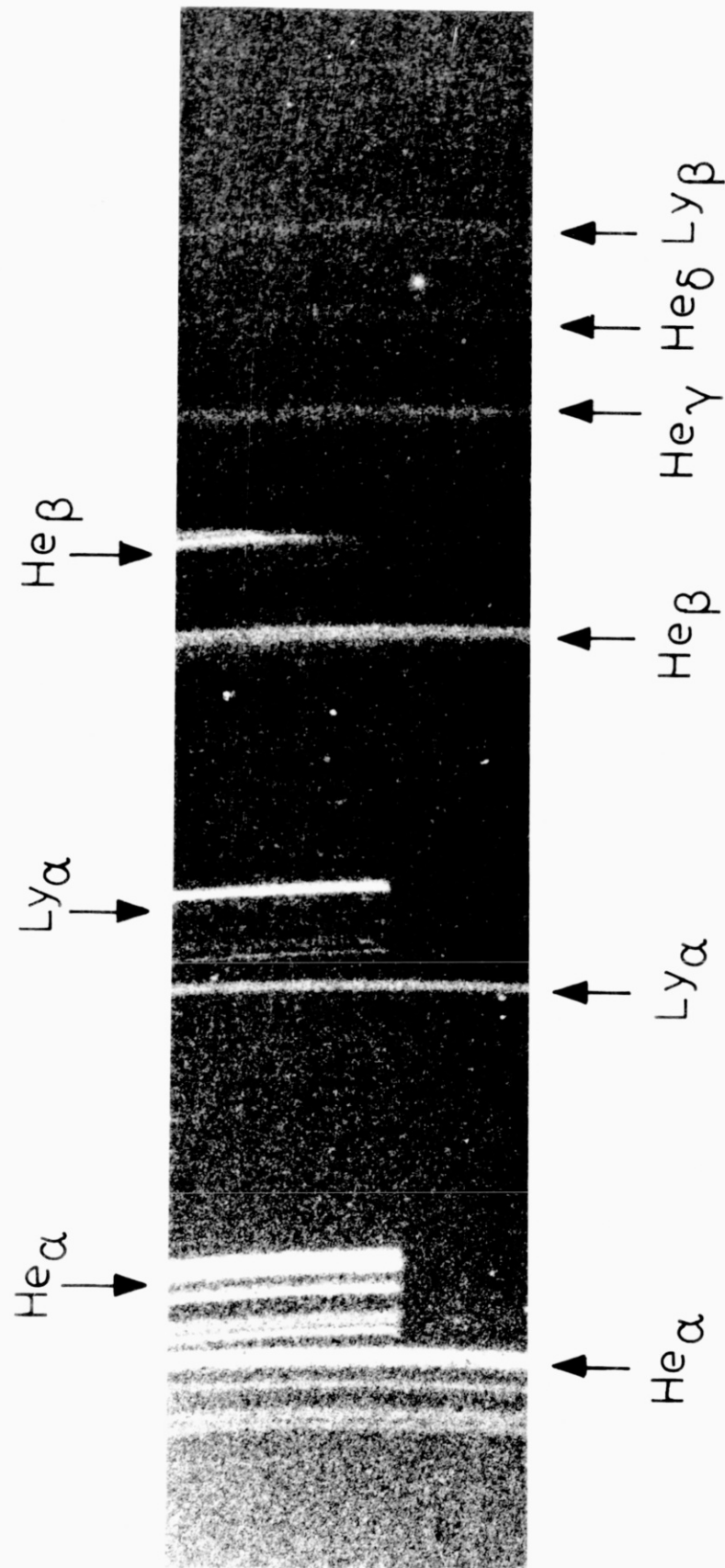


Figure 8.6 Densitometer traces of the streak spectrum presented in figure 8.5, taken at 20 ps and 35 ps after the laser pulse.

TIME INTEGRATED AL SPECTRUM WITHOUT PREPULSE



TIME INTEGRATED AL SPECTRUM WITH PREPULSE

Figure 8.7 A comparison of time integrated X-ray spectra from Al slab targets, with and without a large ASE prepulse.

8.4.2 HIGH DENSITY OBSERVATIONS

The electron density n_e was found for low prepulse shots by comparison of experimental Stark profiles to synthetic profiles generated by the codes RATION and SPECTRA [87]. RATION is a steady state, non-LTE atomic physics code, and while it is not ideal for modelling such short pulse phenomena, at the densities observed ($n_e \approx 2 \times 10^{23} \text{ cm}^{-3}$) ionisation times and excitation times are comparable to or much shorter than the 3 ps laser pulse. Figure 8.8 compares an experimental trace of He_β taken at 15 ps after the laser pulse with a profile generated by SPECTRA for an electron density of $1.4 \times 10^{23} \text{ cm}^{-3}$ and electron temperature of 400 eV.

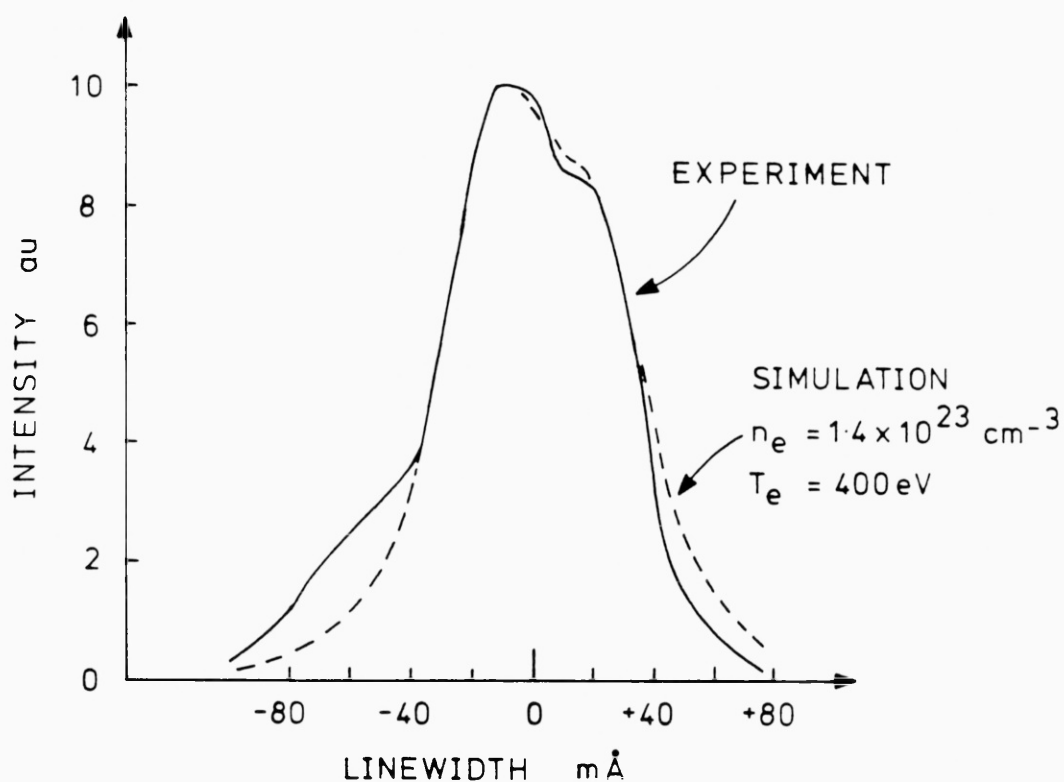


Figure 8.8 A comparison of experimental and synthetic H_β line profiles for $n_e = 1.4 \times 10^{23}$ and $T_e = 400 \text{ eV}$.

The calculations of the spectra used a temperature of 400 eV (see section 8.4.4) and made corrections for opacity broadening [104] which is small as the plasma is only on the order of a micron thick at this time. The fine features of the profiles such as the splitting of He_β are below the finite (15 mÅ) spectral resolution of the X-ray spectrometer, and these details are not well observed. The codes do however give a good fit to the Stark widths and overall profile, indicating a density of $n_e \approx 1.4 \times 10^{23} \text{ cm}^{-3}$, some 10 times higher than critical density for quarter micron KrF light.

The continuum lowering may also be used to find the electron density. Figure 8.9 shows full synthetic spectra generated by RATION and SPECTRA for electron densities of 2×10^{22} , 3.2×10^{23} and $6.4 \times 10^{23} \text{ cm}^{-3}$. There is obviously a significant change in spectral content with increasing density as continuum lowering suppresses higher series members and the remaining lines are Stark broadened. From these simulations the suppression of Al He_γ and $\text{Ly}\beta$ seen experimentally indicates that X-ray emission from densities well above 10^{23} cm^{-3} was seen.

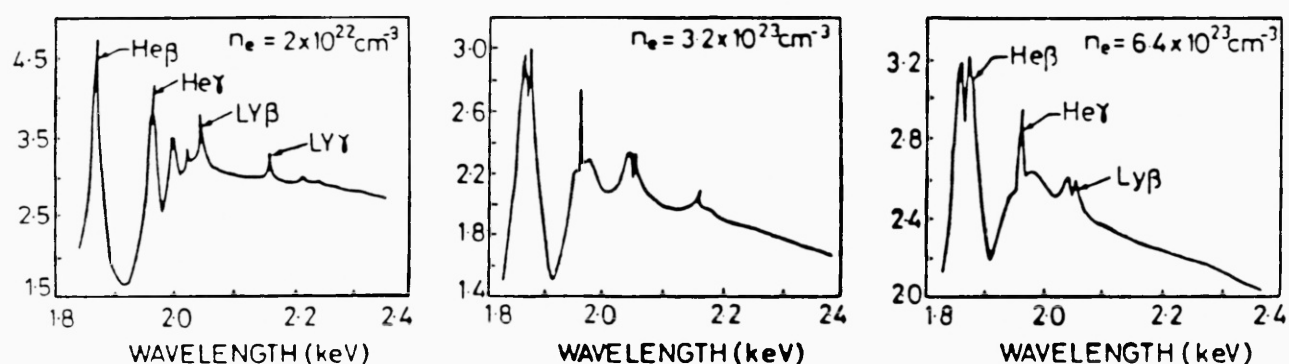


Figure 8.9 Synthetic aluminium spectra for electron densities of 2×10^{22} , 3.2×10^{23} and $6.4 \times 10^{23} \text{ cm}^{-3}$.

A simple analytical model may also be applied to the continuum lowering. A precise theory for this regime does not exist, due to the difficulty of performing the necessary quantum mechanical calculations when the wave functions of different n states begin to mix. However a formulation due to Inglis and Teller [106] can be used to give a rough estimate (within a factor of 2 at best) of the electron density.

For hydrogenic ions the suppression of the $n_s \rightarrow 1$ line requires an electron density :-

$$n_e \approx a_0^{-3} 2^{-15/2} n_s^{-15/2} Z_i^{9/2} \quad (8.1)$$

where a_0 is the Bohr radius, n_s the principle quantum number of the suppressed $n_s \rightarrow 1$ transition and Z_i the perturbed ion charge. This gives a value of $n_e \approx 7 \times 10^{23} \text{ cm}^{-3}$ for the Al $\text{Ly}\beta$ transition, which is again consistent with the more accurate value given by the Stark profile fitting. This quasi-static formulation ignores ion-ion correlations, however these result in an error [107] which is well below the inherent factor of 2 accuracy from this simple order of magnitude argument.

8.4.3 X-RAY SPECTRA WITH A LARGE PREPULSE

The results from 8.4.1 and 8.4.2 may be compared with shots in which the SPRITE laser was not optimised, giving a large 20 ns ASE prepulse containing some 30 % of the short pulse energy. Figure 8.10 shows a typical optical and X-ray diode summed trace for such a shot. The diode response is highly non linear. The X-ray streak spectrum in figure 8.11 is from an Al foil target irradiated with such a pulse, and in this case the prepulse breaks down the target resulting in the short pulse interacting with a large plasma corona. The spectrum is distinctly different from figure 8.5 and much more typical of a critical density Al plasma showing the full H and He like series.

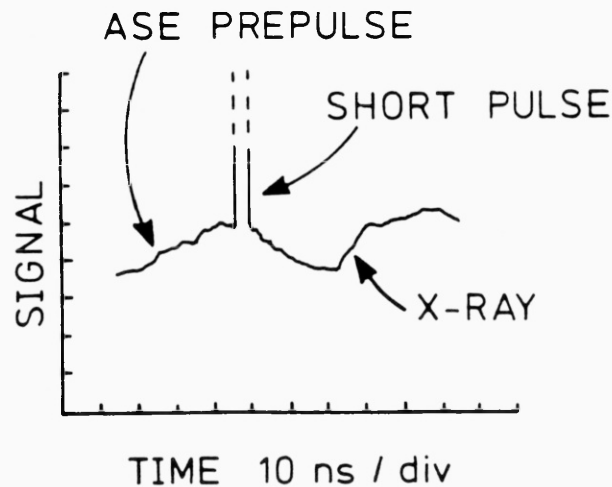


Figure 8.10 Optical diode and X-ray diode summed trace for a SPRITE shot with large ASE prepulse.

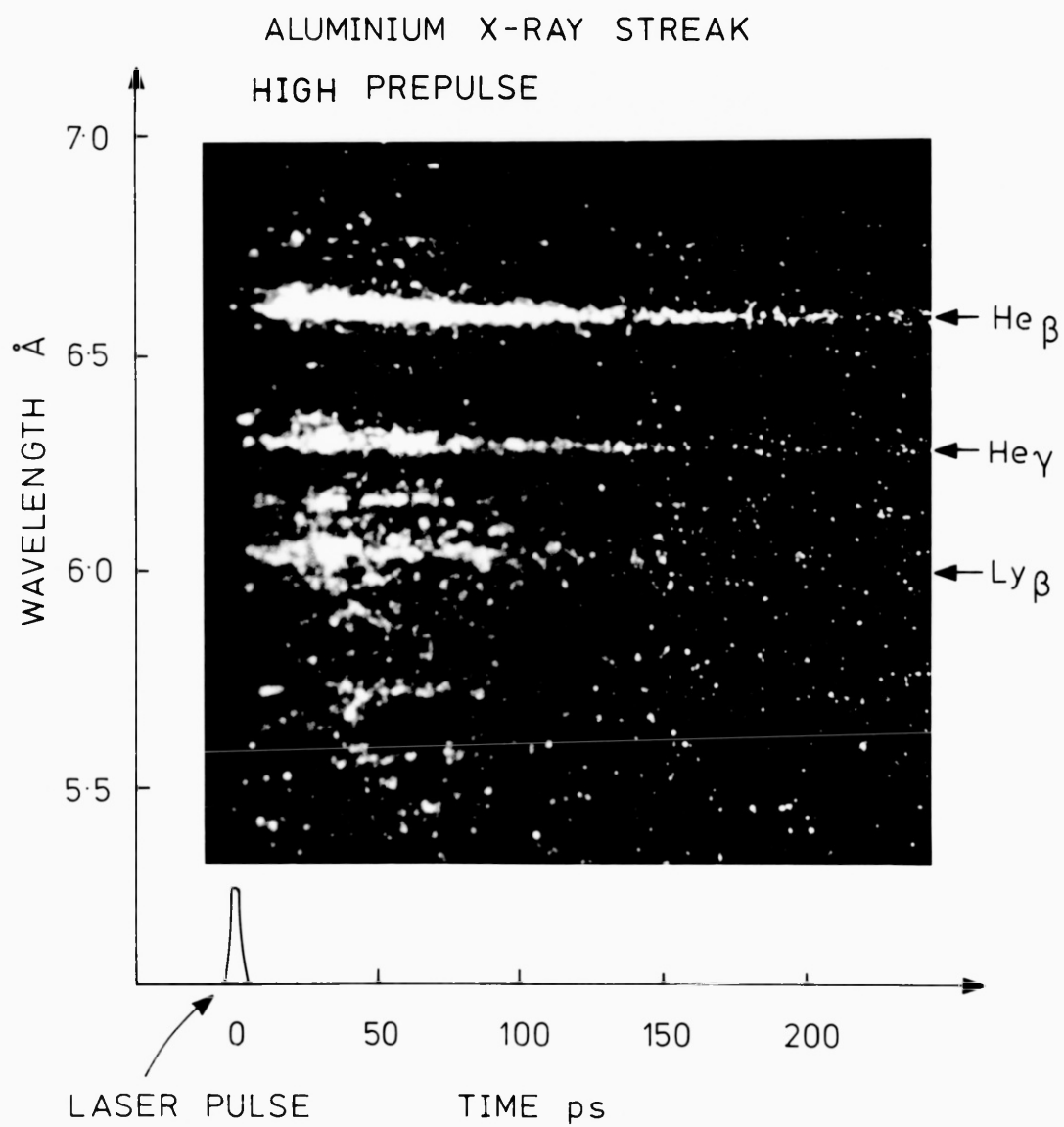


Figure 8.11 X-ray streak spectrum from an Al slab target irradiated with a short pulse superimposed on a large ASE background.

These lines are emitted for several hundred picoseconds without significant change in brightness, indicating a large plasma corona which expands and cools more slowly. Similarly, the time integrated spectra for these shots (figure 8.7) show higher series members and lower intensity satellites, again more typical of lower density plasmas.

The electron density was again found by comparison with synthetic spectra produced by RATION and SPECTRA and good agreement was found for $n_e \approx 3 \times 10^{22} \text{ cm}^{-3}$. This is consistent with X-ray emission from densities close to critical ($1.6 \times 10^{22} \text{ cm}^{-3}$) for $0.248 \text{ } \mu\text{m}$ light.

8.4.4 TEMPERATURE DETERMINATION FROM X-RAY SPECTRA

The electron temperature T_e was found from the ratio $\text{Al Ly}_\alpha / \text{Al He}_\alpha$ in time integrated spectra. Ideally a pair of higher series H and He like lines should have been used as they are optically thinner (having lower oscillator strengths) and more nearly in thermodynamic equilibrium with the electron continuum. This is however impossible to do at such high ($< 10^{23} \text{ cm}^{-3}$) densities as all the series members at wavelengths shorter than Al He_β are suppressed by continuum lowering. Al Ly_α did not fall within the spectral range of the time resolved X-ray spectrometer, and so a time integrated spectrum was used. However the plasma is very short lived, and a reasonable estimate of the temperature can be obtained after the laser pulse as the target is heated by superthermal electrons. Using estimates of self absorption coefficients to correct for opacity [108] a temperature of $400 \pm 50 \text{ eV}$ was found.

8.4.5 XUV SPECTRAL OBSERVATIONS

Time resolved XUV spectra were recorded on a number of shots. Figure 8.12 shows a typical XUV streak from an Al foil target irradiated at $3 \times 10^{16} \text{ Wcm}^{-2}$. The emission is clearly much longer lived than the X-ray signals from similar shots and is composed primarily of higher order Li and Be like transitions. No H or He like Al lines are seen in higher order as this grating is very inefficient below $15 \text{ } \text{\AA}$. The large number of lines in the spectrum results in many line overlaps, and this together with the low resolution ($0.5 \text{ } \text{\AA}$) and low dispersion of the instrument make it impossible to compare experimental line profiles with synthetic spectra to find the density of the emitting plasma.

AL XUV STREAK SPECTRUM

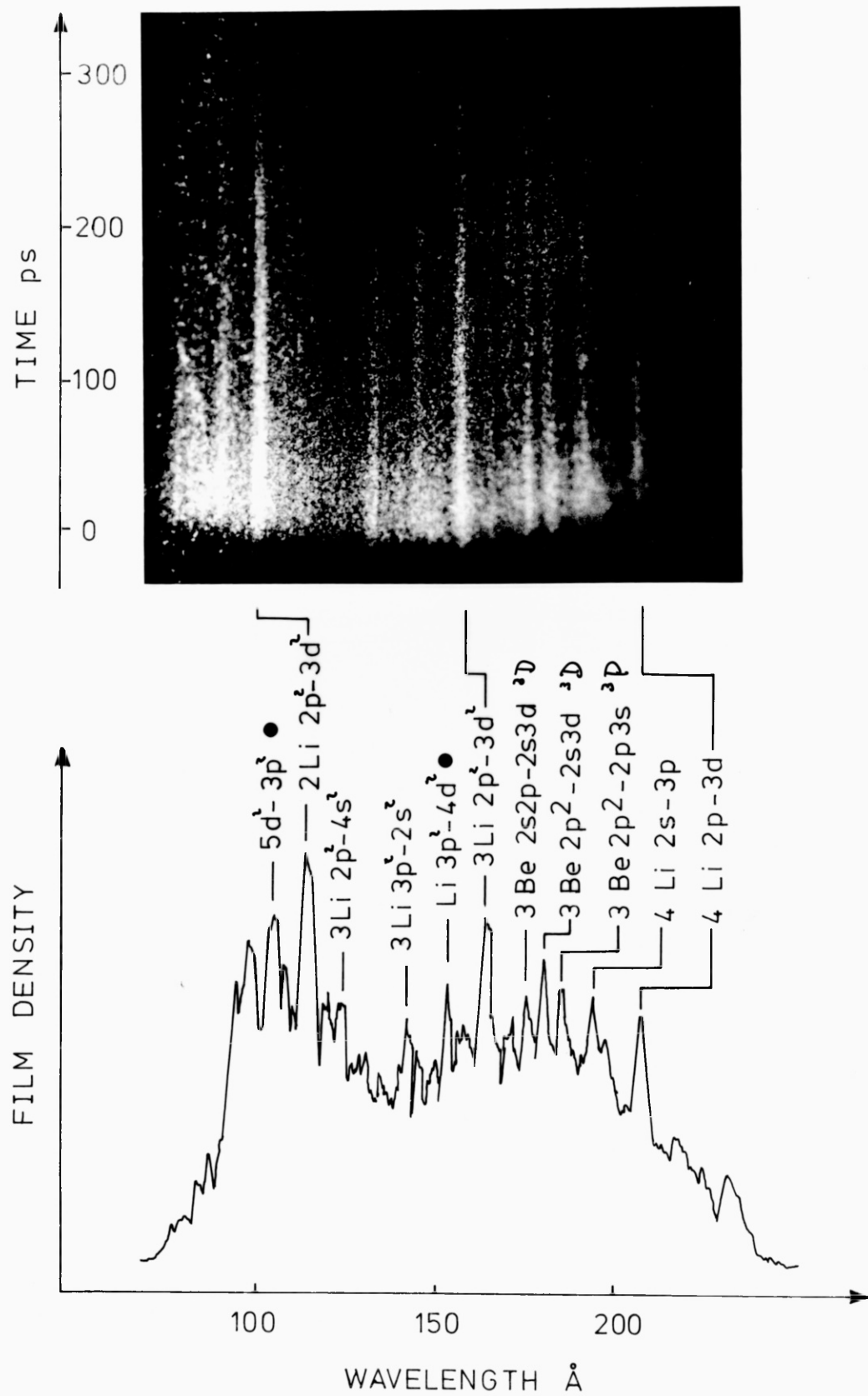


Figure 8.12 XUV streak spectrum from an Al slab target irradiated with a SPRITE short pulse at an intensity of $3 \times 10^{16} \text{ Wcm}^{-2}$.

It is interesting to note that the turn on of the XUV spectrum is very sharp, and that there is no obvious XUV precursor caused by the low ASE prepulse indicating that the ASE does not break down the target. The effect of a large ASE background on the VUV spectrum will be discussed later.

One particularly noteworthy feature is that the Li like Al laser lines 5d-3p and 4d-3p are considerably brighter (by about a factor of 20) than those seen in the transverse spectra from Al fibre X-ray laser targets. A similar effect is also seen in carbon spectra. A lower Z material such as carbon gives a less complex spectrum, as can be seen in figure 8.13 which shows an XUV streak from a CH slab target. The spectrum is composed of high order H and He like resonance lines and is dominated by 3, 4, 5 and 6 Ly α which are noticeably Stark broadened at early time. Again the laser line Balmer H α at 182 Å is brighter by a factor of 8 or so compared to the laser / resonance line ratio from transverse X-ray laser experiments. This spectrum can be compared to carbon spectra from longer pulse (70 ps) X-ray laser experiments (chapter 4, figure 4.13) where the transverse spectrum of a C fibre is again completely dominated by high order resonance lines, but H α is one of the faintest lines in the spectrum, even though there is a 3-2 inversion and significant gain in the axial direction.

It is possible that in the short pulse case the rapid hemispherical expansion of the plasma leads to super-cooling and the formation population inversions at early times when the density is high. ASE with large gain due to the high density could then produce the unusually high laser line intensities relative to resonance lines seen in this experiment. Indeed early experiments on 3-2 inversions in carbon laser plasmas were carried out in planar geometry with point focus irradiation [32] but the plasma would have expanded from a much lower initial density. If this is the case then it is very encouraging for future X-ray laser studies with short pulses. There has been some discussion as to whether this effect could be properly described as super-radiance. Allen and Peters [88] suggest that amplified spontaneous emission (ASE) is more appropriate in this case.

In contrast to figure 8.12, when a large ASE prepulse was present, a faint XUV precursor to the main XUV signal was seen. Figure 8.14 compares XUV streak spectra from Al foil target irradiated by SPRITE pulses, with and without a large ASE background. A faint XUV precursor is visible in the prepulse case, and the most prominent feature appears to be a series of high order Be like Al resonance lines suggesting that the prepulse now breaks down the target but does not strip the ions beyond Al⁹⁺ since no Li like lines are seen.

CH XUV STREAK SPECTRUM

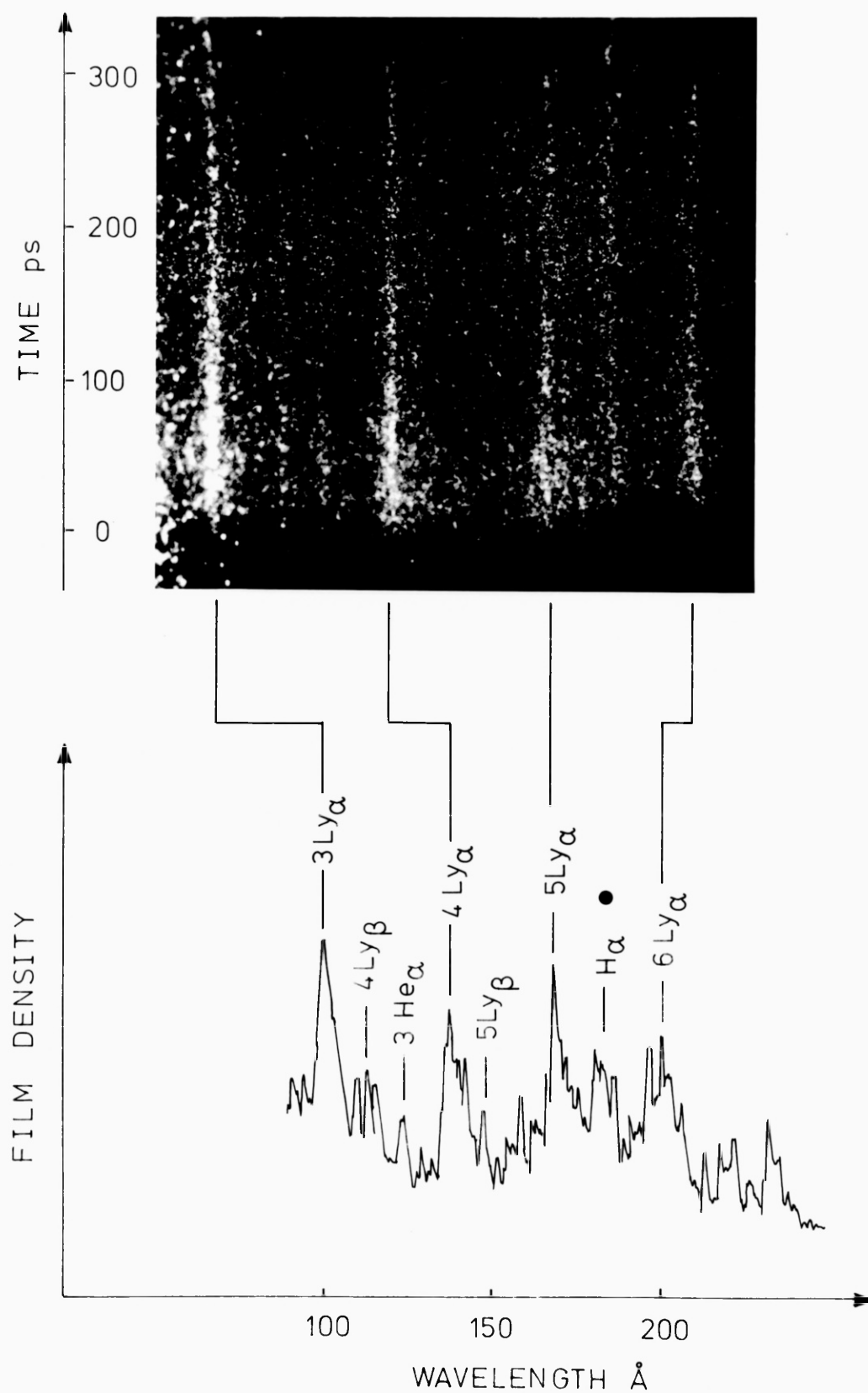


Figure 8.13 XUV streak spectrum from a CH slab target irradiated with a SPRITE short pulse.

VUV STREAK SPECTRA FROM AL
FOIL TARGETS WITH AND WITHOUT
PREPULSE

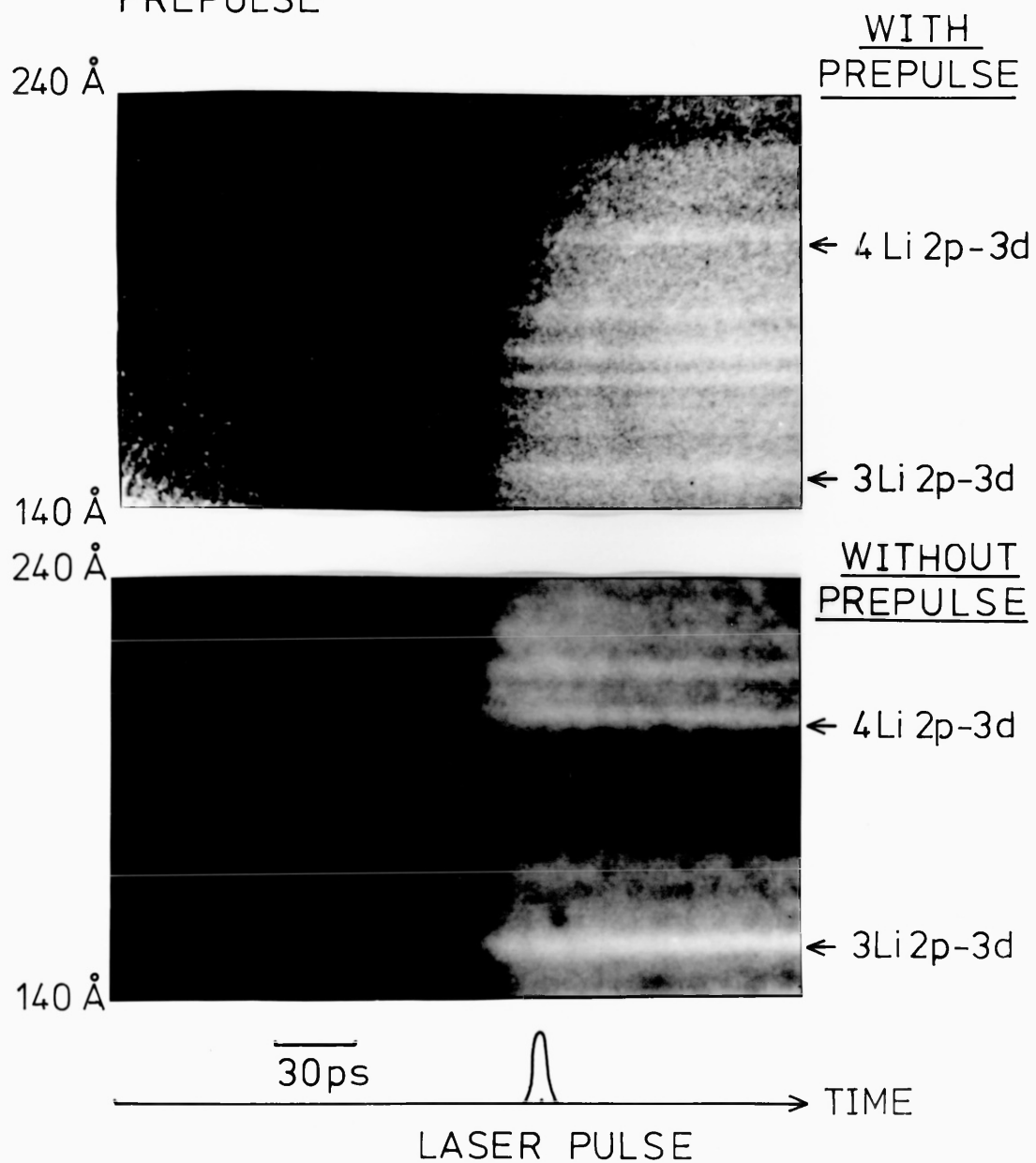


Figure 8.14 A comparison of XUV streaks from AL slab targets irradiated with SPRITE short pulses, with and without a large ASE background.

8.4.6 Ti K_{α} STUDIES AND HOT ELECTRON PRODUCTION

Titanium foil targets coated with various thicknesses of CH were used to investigate the level of hot electron production in shots with and without a large ASE prepulse. Ti K_{α} has an excitation energy of 4.51 KeV and so any K_{α} signal seen would be due to direct excitation by > 4.51 KeV electrons rather than the 400 eV thermal electrons in the bulk of the plasma. The K_{α} signal was measured with time integrated X-ray crystal spectrometers, and with a prepulse of $< 10^{16} \text{ Wcm}^{-2}$ only K_{α} emission was seen from bare Ti targets. Figure 8.15 shows K_{α} spectra for bare Ti foils and foils coated with $0.5 \mu\text{m}$ of CH. For a $0.5 \mu\text{m}$ overcoat of CH a similar but smaller K_{α} signal was seen, but no K_{α} signal was seen with $\geq 2 \mu\text{m}$ of CH overcoat. From the K_{α} intensity on these shots it was found that $\approx 20\%$ of the incident laser energy was deposited in > 4.51 KeV hot electrons. For this low prepulse case the ASE irradiance is below the breakdown threshold of the target, and this is supported by the prepulse failing to burn through $0.5 \mu\text{m}$ of plastic.

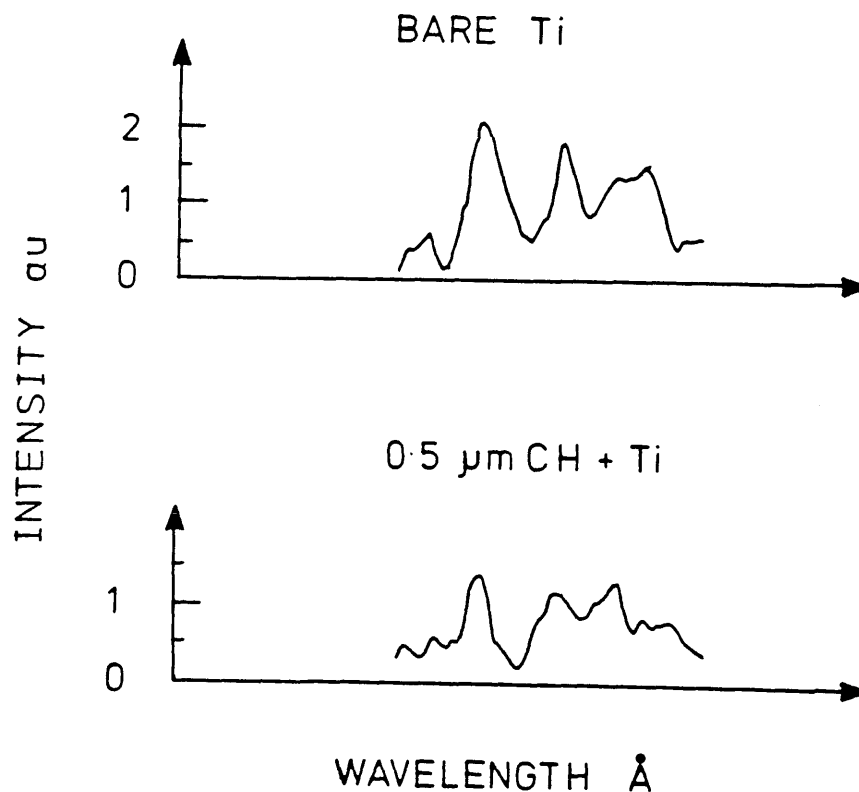


Figure 8.15 K_{α} spectra from bare and CH coated Ti foil targets irradiated with SPRITE short pulses.

In contrast figure 8.16 shows a spectrum from a Ti foil target with 0.5 μm plastic overcoat irradiated with a short pulse superimposed on a large ASE background with a prepulse irradiance well above 10^{11} Wcm^{-2} on target. In this case a strong Ti He_{α} signal and associated satellites are seen in addition to a large K_{α} signal. The high prepulse irradiance broke down the target surface to form a large underdense plasma. The short pulse coupling efficiency was much higher due to inverse Bremsstrahlung absorption in the large underdense region. This gave a much hotter but lower density plasma than in the low prepulse case and allowed the formation of He like Ti which implies a plasma temperature of at least 1.43 KeV.

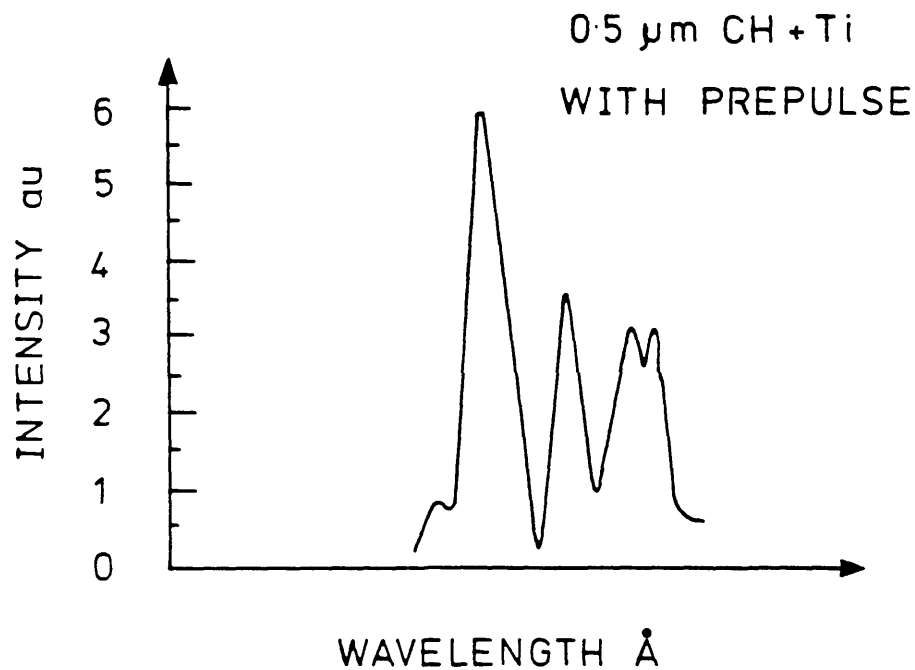


Figure 8.16 K_{α} spectrum from a CH coated Ti foil target, irradiated with a short pulse superimposed on a large ASE background.

8.5 NUMERICAL SIMULATIONS

Simulations of short pulse KrF interactions with and without a prepulse were carried out to help with the interpretation of experimental results. The low prepulse simulations were run for $\lambda = 0.248 \mu\text{m}$, 0.55 J, 3.5 ps (FWHM) gaussian pulses to model the experimental shot in figure 8.5. A large number of 1D hydrodynamic simulations were run with the MEDUSA [89] Lagrangian hydrocode for a range of flux limiters and absorbed energies.

The laser energy was assumed to be absorbed by inverse Bremsstrahlung up to the critical surface and a fraction F_r of the remaining energy dumped at the critical surface to simulate resonance absorption. The fraction of resonantly absorbed energy going into hot electron production was assumed to be 90 % and a semi-empirical hot electron spectrum due to Giovanelli [92] was used to model their spread in energy.

The best agreement with experimental observations of electron density N_e and electron temperature T_e were found for $F_r = 20 \%$ and a flux limiter of 0.1 times the classical free streaming limit. The value of F_r found in simulations was in good agreement with the absorption fraction obtained from K_α measurements [section 8.4.6], and this indicated that very little energy was absorbed by inverse Bremsstrahlung. This is an expected result as there is insufficient time for a large sub-critical plasma to form during the 3 ps laser pulse. Figure 8.17 shows MEDUSA outputs of n_e , T_e and Z_i the average ionisation stage for an Al slab target during and after the laser pulse. The code predicts higher densities during the pulse than are seen experimentally and this is due to the 15 ps resolution of the X-ray spectrometer giving a temporal smearing of the two regimes in figure 8.17.

Shots with a large ASE prepulse were also simulated, with a 20 ns Gaussian containing 30 % of the short pulse energy being superimposed on a 3 ps, 1.7 J pulse. A flux limiter of 0.03 times the classical free streaming value was used as the long prepulse allows the formation of a large under-dense plasma before the arrival of the short pulse. During the prepulse several microns of target are ablated and the electron temperature falls off sharply above critical density (See figure 8.18). This is in good agreement with experimental observations of X-ray emission from near critical density on shots with a large prepulse

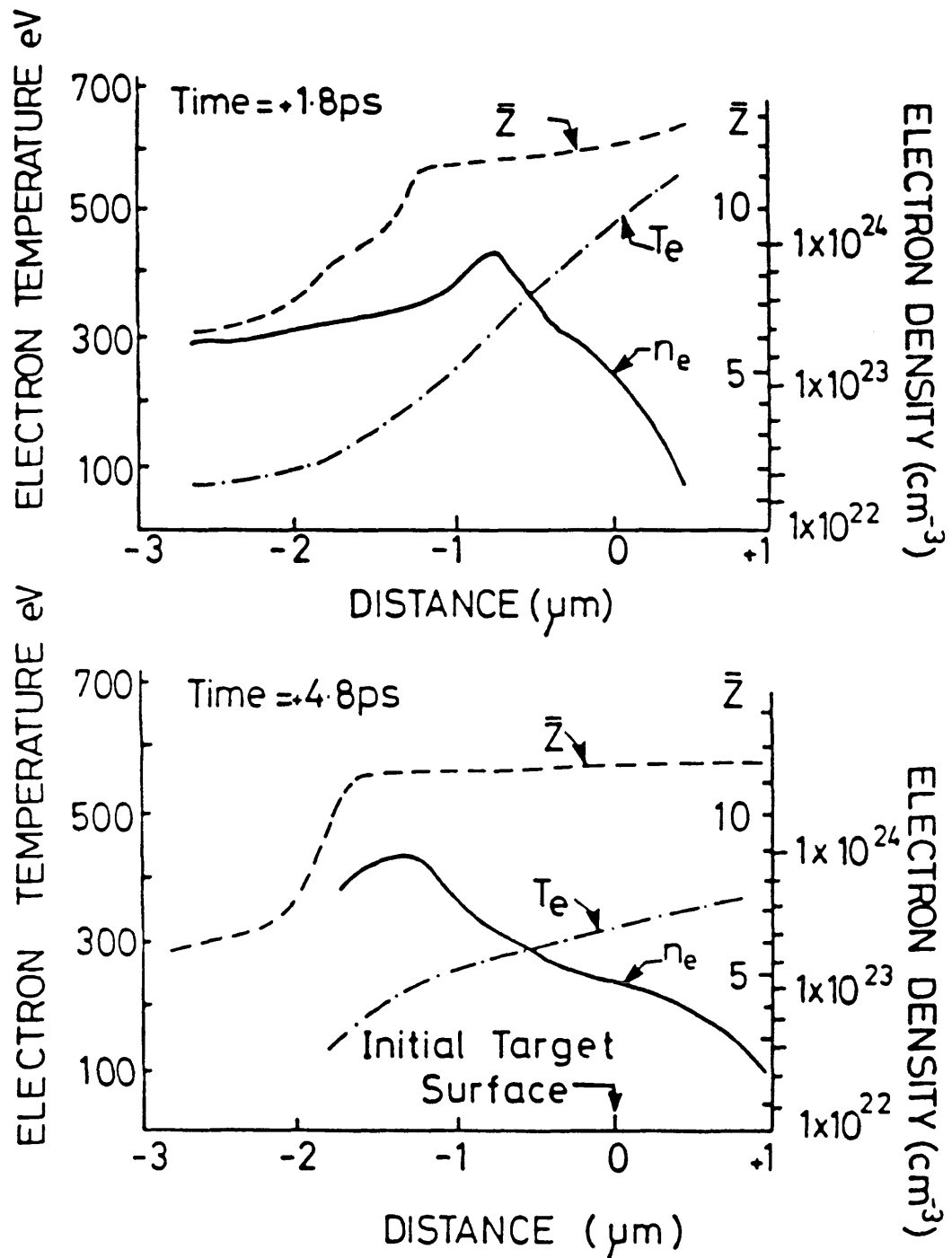


Figure 8.17 MEDUSA 1D hydrocode simulation of n_e , T_e and Z_i for a SPRITE short pulse interacting with an Al slab target.

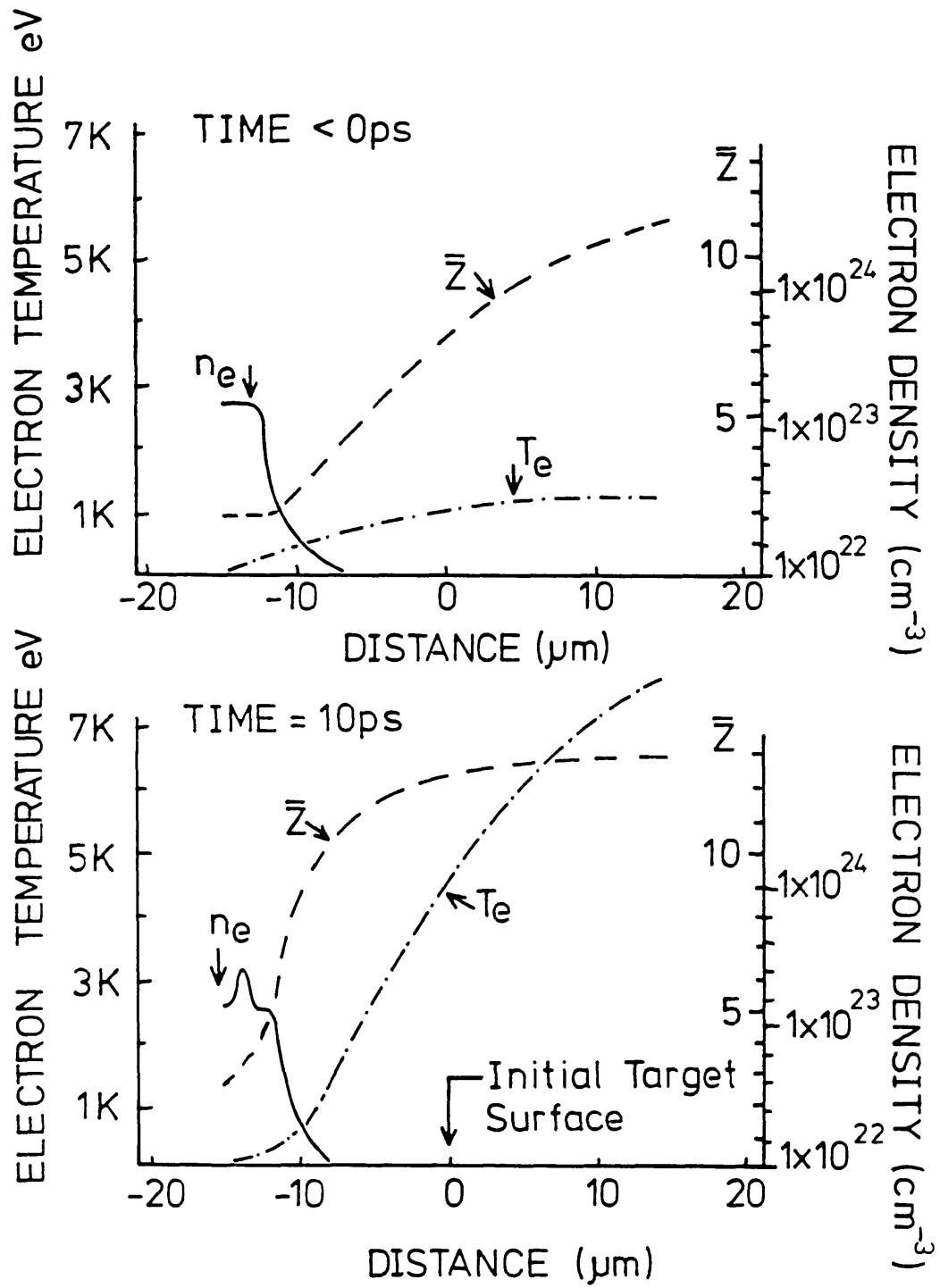


Figure 8.18 MEDUSA 1D simulation of a SPRITE short pulse superimposed on a large ASE background, interacting with an Al slab target.

A limited number of 2D Fokker–Planck simulations were also run using a recently developed code [109]. Figure 8.19 shows a simulation for a 3 ps Gaussian pulse irradiating half a 20 μm (FWHM) focal spot at an intensity of $3 \times 10^{16} \text{ Wcm}^{-2}$. (Only half the focal spot is modelled as the interaction is symmetrical and this saves on computer time). Contours are isotherms (T_e in electron volts) and arrows indicate the direction and magnitude of heat flow.

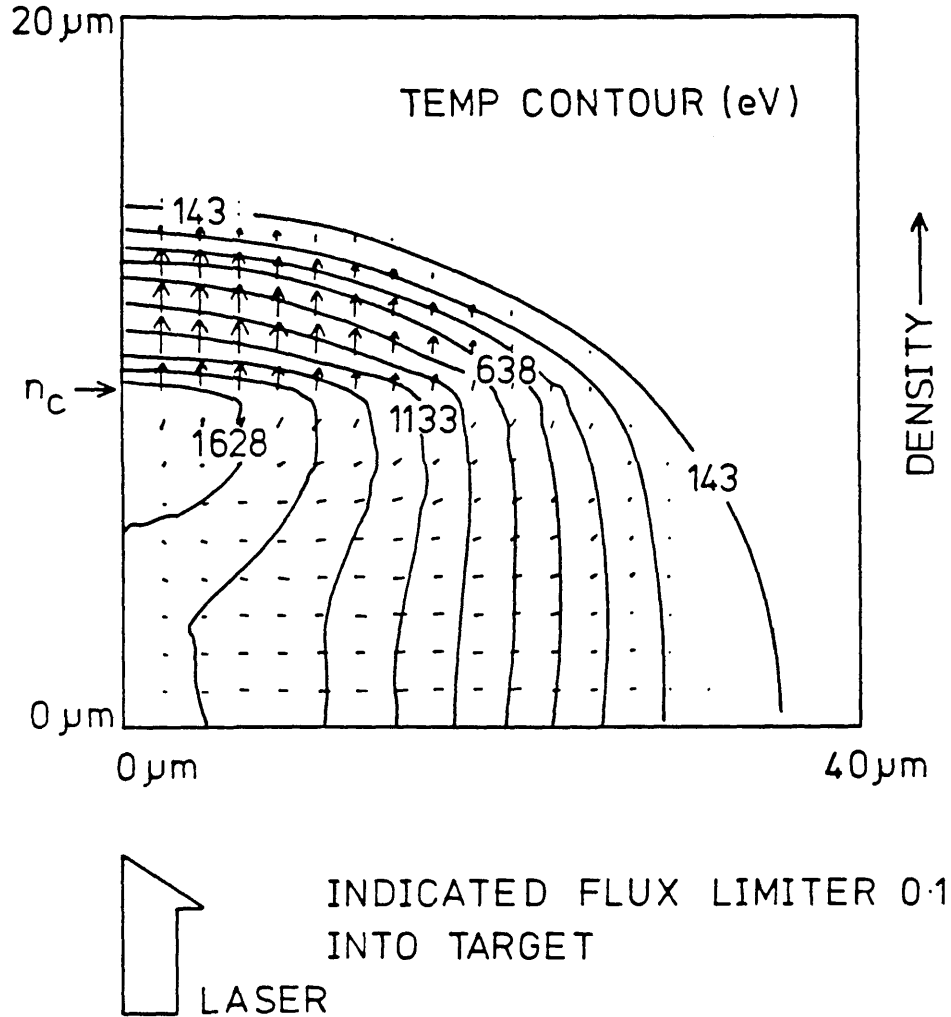


Figure 8.19 2D Fokker–Planck simulation of a 3 ps Gaussian short pulse interacting with an Al slab target. Contours are T_e isotherms, arrows indicate the direction and magnitude of heat flow.

It can be seen that very little energy is absorbed below critical density, and the majority of the energy absorbed at the critical surface is transported directly into the the target. The lateral heat flow that does take place is well out into the plasma corona, below critical density where the density is low and transport of energy away from the focal spot is insignificant.

This simulation explicitly gives a value for the flux limiter of 0.1 times classical and predicts electron densities of $\approx 2 \times 10^{23} \text{ cm}^{-3}$ and temperatures of $\approx 400 \text{ eV}$ at the center of the focal spot, immediately after the short pulse interaction. These are consistent with experimental observations and 1D MEDUSA simulations. The good agreement of both 1D hydro and 2D Fokker–Planck codes with experimental observations indicates that the results are well understood, and that lateral heat transport which is not included in the 1D code is relatively unimportant on the picosecond timescale.

8.6 CONCLUSION

The first single shot picosecond KrF interaction study with solid targets was carried out. With low ($< 10^{11} \text{ Wcm}^{-2}$) ASE prepulse irradiances on target, X-ray Stark profiles and continuum lowering indicated that emission was seen from fully stripped aluminium plasmas at electron densities of well above 10^{23} cm^{-3} , approaching solid density. X-ray line ratio measurements gave a temperature of $400 \pm 50 \text{ eV}$ for these plasmas. 1D hydro and 2D Fokker–Planck code simulations were in good agreement with experimental values of n_e and T_e as well as with the 20 % coupling efficiency to hot (4.51 KeV) electrons seen in K_{α} emission studies.

In contrast, with large prepulse levels, both experiments and simulations indicated plasmas at much lower densities close to critical, and with temperatures in the few KeV range.

This work is encouraging for the scaling of recombination lasers to shorter wavelengths. In particular an H like Al water window laser would require a fully ionised plasma expanding from near solid density to achieve significant gain. With respect to this, unusually intense emission of carbon H_{α} at $182 \text{ \AA}$ from rapidly cooling high density plasmas was observed, and may have been due to ASE.

CHAPTER 9. SUMMARY

This thesis has described work carried out by the author on several XUV laser schemes, and experiments on related laser plasma interactions. Work in this field was recently stimulated by demonstrations of XUV gain at 206 Å and 209 Å in laser produced Se^{24+} plasmas at the Lawrence Livermore Laboratory. In the three years covered by this thesis, the shortest wavelength at which gain has been observed has fallen to just below 50 Å.

Much of the recent success in the field has been due to improvements in diagnostics and experimental techniques. Chapter 3 outlines the construction and calibration of a novel flat field, time resolved XUV spectrometer for use in X-ray laser and laser plasma interaction studies. One of the major advantages of these instruments over time integrated diagnostics is their ability to distinguish a bright but short lived ASE signal from the longer lived background plasma emissions. Chapter 4 describes the use of spectrometers of this type to demonstrate significant XUV gains on the order of 4 cm^{-1} at 182 Å in adiabatically cooled, recombining carbon plasmas. These plasmas were produced by laser irradiation of thin fibre targets using a novel off axis lens and mirror facility at the Central Laser Facility of the Rutherford Appleton Laboratory.

Initial disagreements between experiments and detailed numerical simulations on the magnitude and time evolution of gain described in chapter 4 have resulted in both improvements to diagnostic techniques and the inclusion of important additional physics in computer codes. In consequence experimental observations and numerical simulations now show good convergence for several recombination XUV laser schemes.

The hydrogenic XUV laser scheme was then scaled to shorter wavelength for the first time by moving to a higher Z lasant ion. Chapter 5 describes what is currently the most completely diagnosed and fully modelled X-ray laser experiment to have been carried out. This was also the first experiment to show XUV gain below 100 Å and the detailed analysis has since produced valuable data on the optimisation of plasma conditions for gain in recombining plasmas. Three plasma energy regimes were investigated, and gain coefficients of $-1.8 \pm 0.2 \text{ cm}^{-1}$, $4.4 \pm 1 \text{ cm}^{-1}$ and $2.0 \pm 1 \text{ cm}^{-1}$ found for the H_α line at 81 Å for absorbed plasma energies of 5, 9 and 16 J cm⁻¹ respectively.

In this study simultaneous axial and transverse time resolved spectroscopy was used for the first time at RAL to diagnose an X-ray laser experiment, and resulted in much more accurate observations of the time evolution of gain than had previously been attained.

Hydrogenic recombination schemes have been highly successful in producing XUV gain in the 200 – 50 Å spectral range. However lithium like ions are also possible XUV amplifiers and have significant energetic advantages over hydrogen like ions. Chapter 6 reports on a lithium like aluminium XUV laser experiment which demonstrated gains of the order 2 cm^{-1} on 5d–3p, 5f–3d, 4d–3p and 4f–3d transitions at 103.8, 105.7, 150.7 and 154.6 Å. Fibre targets were used to enhance the escape of the lower laser level to ground state decay radiation by increasing Doppler decoupling in the radial direction. As a result, gains a factor of two higher than obtained by previous workers using slab targets were observed.

The scaling of recombination schemes to shorter wavelengths by moving to higher Z lasing ions imposes serious limits on the initial plasma density needed for optimum gain. This high density requirement would be somewhat relaxed if the usual adiabatic cooling of the expanding plasma could be enhanced in some way. In addition, working X-ray laser schemes could be made more energy efficient as population inversions would be formed at higher densities. Chapter 7 examines radiative cooling by high Z dopants for the hydrogenic fluorine XUV laser scheme.

A range of solid fluoride slab targets were irradiated in point focus, and the cooling phase of the plasma diagnosed with time resolved X-ray spectroscopy. The electron temperature was found by comparing experimental line ratios to atomic physics code predictions. Similar cooling rates were found for targets containing F, Na and Mg as the highest atomic number element, reflecting their similar atomic numbers. However a 2–3 fold enhancement in cooling rate due to increased radiation loss from the plasma was seen for targets containing K and Ca dopants.

The scaling of recombination X-ray lasers to still shorter wavelengths will require laser plasmas produced at high initial density. For example a hydrogen like aluminium water window recombination laser would need a plasma expanding from near solid density. Such plasmas are difficult to make in a geometry suitable for producing gain with pump laser pulses longer than 20 ps, because the leading edge of the pulse ablates a large underdense plasma.

In this "long pulse" case the majority of the pump energy is coupled to the plasma by inverse Bremsstrahlung at some distance from the target surface. Chapter 8 reports on a laser plasma interaction study in which the pulse length (≈ 3 ps) was short compared to the hydrodynamic expansion time of the plasma. This resulted in energy being dumped close to the initial target surface, giving hot (≈ 400 eV) hydrogen like plasmas at near solid density ($\approx 10^{23}$ cm $^{-3}$).

In summary this thesis has reported on a novel diagnostic developed in part by the author. This instrument was used by the author on a number of highly successful X-ray laser experiments, and several more basic laser plasma interaction studies of relevance to X-ray laser research. These experiments have contributed both to the field of X-ray lasers, and more generally in the area of high power laser plasma interactions. Several refereed publications have resulted from this work [8,9,10,50]. In addition much of this work has also appeared in published conference proceedings [18,19,91 and 117 to 126] and unrefered reports for the Rutherford Appleton Laboratory [112 to 116]. Two further publications on the axial to transverse carbon X-ray laser and fluorine radiative cooling experiments are presently in preparation.

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APPENDIX I. COLLABORATORS AND CO-WORKERS

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+ Queen's University Belfast

o Rutherford Appleton Laboratory

• University of York

△ LSAI, Orsay, France

▽ Electrotechnical Laboratory, Ibaraki, Japan

Chapter 3. Spectrometer calibration

Experimental. G. Kiehn*, R.A. Smith*, T. Garvey*, O. Willi*

Analysis. G. Kiehn*, R.A. Smith*

Chapter 4. Axial carbon laser experiment

Experimental. XUV Spectroscopy G. Kiehn*, R.A. Smith*

Pinhole cameras C. Regan⁺

Supervision. O. Willi*, C. Lewis⁺, M. Key^o

Analysis. G. Kiehn*

Simulations. S. Rose^o, G. Pert*

Chapter 4. Axial / transverse carbon experiment

Experimental. Axial XUV Spectroscopy R.A. Smith*

Transverse XUV Spectroscopy D. O'Neal⁺

Pinhole cameras / Calorimetry C. Regan⁺

Analysis. R.A. Smith*

Simulations. S. Rose^o

Chapter 5. Fluorine laser experiment

Experimental. Axial XUV Spectroscopy G. Kiehn*

Transverse XUV Spectroscopy R.A. Smith*

Pinhole cameras / Temp Monitor C. Regan⁺

Calorimetry / X-Ray Spectroscopy T. Tommi[▽], M. Grande^o

Supervision. O. Willi*, C. Lewis⁺, M. Key⁰
 Analysis. Plasma regimes / Transverse spectra / Axial transverse ratios
 Time evolution of gain R.A. Smith*
 Axial spectra G. Kiehn*
 Pinhole data C. Regan⁺
 Simulations. S. Rose⁰, G. Pert*

Chapter 6. Lithium like aluminium experiment

Experimental. Axial XUV Spectroscopy C. Regan⁺
 Transverse XUV Spectroscopy R.A. Smith*
 Pinhole cameras D. O'Neal⁺
 Supervision. O. Willi*, C. Lewis⁺, M. Key⁰
 Analysis. Axial spectra / Data set selection and
 Time evolution of gain C. Regan⁺
 Transverse spectra R.A. Smith*
 Pinhole Data D. O'Neal⁺
 Simulations. A. Klisnick[△]

Chapter 7. Radiative cooling

Experimental. 2400 XUV Spectroscopy R.A. Smith*
 1200 XUV Spectroscopy J. Edwards*
 Support G. Kiehn*, O. Willi*
 Analysis. R.A. Smith*
 Simulations. R.A. Smith*

Chapter 8. Picosecond Spectroscopy

Experimental. X-Ray Spectroscopy G. Kiehn*, V. Barrow*
 XUV Spectroscopy R.A. Smith*, J. Edwards*
 Supervision. O. Willi*
 Analysis. G. Kiehn*, V. Barrow*, J. Edwards*
 Simulations. J. Edwards*, V. Barrow*

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