

Fluorescence scattering in a cool Z-pinch plasma and
in the HBTX1A reverse-field pinch.

by

Denise Anne Evans, B.Sc.

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Blackett Laboratory,
Imperial College,
London SW7 2BZ,
United Kingdom

ABSTRACT.

This thesis describes laser induced fluorescence studies of two separate plasma devices. First, the recombination phase of a z-pinch plasma run in hydrogen was studied. Previous fluorescence scattering work performed on this plasma had revealed discrepancies between the observed fluorescence signal and the predictions of a 20 level collisional radiative (CR) model of the plasma employing a widely accepted set of electron-atom collision rates.

An alternative set of rates which were typically a factor of 3 lower than the previous set was incorporated into the CR code. Error bars corresponding to the uncertainty in the diagnosed electron temperature were also incorporated. As a result, some of the discrepancies between theory and experiment were reduced. However, the pre-laser density of the first excited state of hydrogen and the fluorescence plateau enhancement were both still much higher than predicted by the CR code. Inclusion of rates into the atomic levels by processes involving H^- and H_2^+ ions made little difference to the predicted values.

Experimental studies of the effect of H^- ions are reported first. A high power Balmer- α fluorescence scattering experiment was performed and the temporal behaviour of the plateau observed. A second experiment was carried out using a ruby laser to photodetach H^- ions and photoionize atomic levels with $n \geq 3$, the repopulation of the $n=3$ and 4 levels being studied. From the results obtained, it was concluded that either the actual H^- ion density was greater than calculated from LTE or the mutual neutralization rate into the $n=2$ level was larger than calculated for the delay times of less than 50 μs . Also, the $n=4$ and $n=3$ levels (for the later delay times) were concluded not to be significantly populated by processes involving H^- ions.

Balmer- α fluorescence scattering was also carried out in the recombination phase of a similar z-pinch device to see if the discrepancies were peculiar to the particular plasma studied previously. The observed fluorescence largely demonstrated the same anomalous behaviour.

The second plasma studied was the HBTX1A reverse field pinch at the UKAEA Culham Laboratories. Laser induced fluorescence at the Balmer- α wavelength was used to measure the neutral deuterium density with a spatial resolution of 4 cm. Balmer- β fluorescence yielded results with higher fluorescence enhancements and a spatial resolution of down to 2 cm. The neutral density profiles so obtained are discussed.

A feasibility study was also carried out to investigate the use of fluorescence scattering to measure the density of highly ionized impurity species. Two transitions involving the Ti XVII ion were considered. However the high saturation power and long time required to attain saturation and the small enhancements available are major obstacles to the use of fluorescence scattering for this purpose.

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Chapter One

INTRODUCTION

1.1 Overview of This Thesis.

This thesis describes the use of the technique of laser induced fluorescence scattering in: a) the recombination phase of a low temperature ($0.2 < T_e < 1.0$ eV) hydrogen z-pinch plasma device at Imperial College and b) the high temperature ($T_e = 100$ eV) reverse field pinch HBTX1A at the UKAEA Culham Laboratory. All the fluorescence measurements were made at the Balmer- α and/or Balmer- β wavelengths. Chapters 2 to 4 inclusive concern themselves with the low temperature work and Chapters 5 to 7 consider the high temperature work.

This first chapter briefly describes the technique of fluorescence scattering, its interpretation and sensitivity to the plasma conditions. Previous work involving fluorescence spectroscopy on low temperature plasmas is also reviewed.

Chapter 2 discusses observed anomalies associated with Balmer- α and Balmer- β fluorescence scattering in the recombination phase of one particular z-pinch plasma device studied at Imperial College. Various explanations for the anomalies observed are also discussed. Chapter 3 describes a set of Balmer- α fluorescence scattering experiments in a plasma device similar to that discussed in Chapter 2. Chapter 4 describes two experiments performed in the same z-pinch device discussed in Chapter 2, both of which investigated the influence of H^- ions on the observed fluorescence signal. The first involved observing the Balmer- α fluorescence signal induced by a laser beam much more intense than used previously, and the second studied the behaviour of the Balmer- α and - β emission during

irradiation by a ruby laser.

Chapters 5 and 6 describe a set of experiments performed at the UKAEA Culham Laboratory on the HBTX1A reverse field pinch. A spatial distribution of the neutral deuterium density was obtained by Balmer- α and - β fluorescence scattering and the results compared to neutral density profiles generated from a transport code. Chapter 7 discusses the possibility of using the fluorescence scattering technique to measure the density of impurities in fusion-type plasmas (particular consideration being given to highly ionized species). A case study for Ti XVII is presented for typical operating conditions for the DITE tokamak.

Finally, Chapter 8 reviews the work described in this thesis and presents several suggestions for further work.

1.1.1 Units and Symbols. The SI system of units has been used throughout this thesis. Where other units are used, this is indicated in the relevant parts of the text.

The main exceptions to the use of SI units are as follows. Particle densities (n or N) are usually quoted in units of cm^{-3} , as is usual in texts on plasma physics, and wavelengths are quoted in Angstroms ($\text{\AA}$). The electron temperature kT_e is in Joules, but the symbol T_e is more commonly used in the text and has units of eV. Rate coefficients (C_{ij}) are in units of cm^3s^{-1} , whereas rates ($K_{ij}=n_i C_{ij}$) have units of s^{-1} (per atom/molecule). Densities and rates usually have subscripts to indicate the relevant levels or transitions. For example, n_e is usually taken to represent the electron density and C_{ij} a particular rate coefficient from the atomic level i to the level j . Symbols such as h, λ and c (speed of light) have their usual meanings except if specified otherwise. In tables the symbol E_x is taken to denote 10^x .

1.2 The Technique of Laser Induced Fluorescence Scattering.

1.2.1 Introduction. The technique of laser induced fluorescence scattering is widely used in many areas of physics. In relation to the work reported in this thesis, it may be described as follows.

An atomic transition is chosen for investigation, and the laser tuned to the corresponding wavelength. The laser must have sufficient power to saturate the transition i.e. to populate the upper and lower levels in the ratio of their statistical weights:

$$\frac{n_j}{n_i} = f_s (g_j / g_i) \quad (1-1)$$

where $n_{i,j}$ are the lower and upper levels of the pumped transition, $g_{i,j}$ their respective statistical weights and f_s is the saturation factor, $0 < f_s < 1$. Observation of the induced fluorescence signal at the same wavelength or at the wavelength of another collisionally coupled transition may give information on the relevant populating and depopulating rates between levels, and may even give the level populations themselves.

An important point to note is that once saturation has been achieved, the signal detected is independent of the laser power, and any increase in the laser power above this does not give an increase in the enhancement seen. This helps in the interpretation of results since detailed knowledge of the absolute laser pump intensity is not required.

For low temperature (< 1 eV) plasmas such as those discussed in Chapters 2,3 and 4, the ground state is effectively decoupled from the other levels. For higher temperature plasmas, such as the fusion-type plasmas discussed in Chapters 5,6 and 7, the ground state is

collisionally coupled to the first excited state. In this case, the fluorescence scattering signal in the visible region of the spectrum (from the excited states) may be related to the absolute ground state density. The application of this technique as a powerful diagnostic for spatially resolved neutral density measurements is discussed in more detail in Chapter 5.

1.2.2 Saturation Intensity. From Gohil (1982) the laser intensity I_λ (measured in $\text{Wcm}^{-2}\text{\AA}^{-1}$) into 4π required to saturate a particular transition with upper state j and lower state i is given by:

$$I_\lambda = \frac{8\pi hc^2}{\lambda^5} \frac{1}{A_{ji}(1-f_s)} \left\{ (K_{ji} + K_j + A_{ji}) - (g_i/g_j) K_{ij} \right\} \quad (1-2)$$

where A_{ji} is the spontaneous radiative decay rate (s^{-1}), K_{ji}/K_{ij} are the collisional de-excitation and excitation rates per atom between the levels and K_j is the total collision rate out of the level j to higher levels.

Figure 1.1 shows the laser intensities required to saturate the Balmer- α and - β transitions (from Equation 1-2) as a function of electron density for an electron temperature of 1 eV and saturation factor $f_s=0.5$, using the set of rate coefficients given by Johnson (1972). Note that the Balmer- β transition requires a larger intensity to produce the same degree of saturation than the Balmer- α transition. This is due to the lower wavelength of the transition, the smaller radiative rate A_{42} and the higher collision rates associated with the higher lying levels. The range of intensities shown in Figure 1.1 are well within the present capabilities of conventional flashlamp pumped dye lasers (powers of up to $1 \text{ MWcm}^{-2}\text{\AA}^{-1}$ have been achieved with an oscillator/amplifier arrangement by the author). Note that in the case of the Lyman- α transition ($\lambda=1215 \text{ \AA}$), not only is much more power required to saturate the transition (by

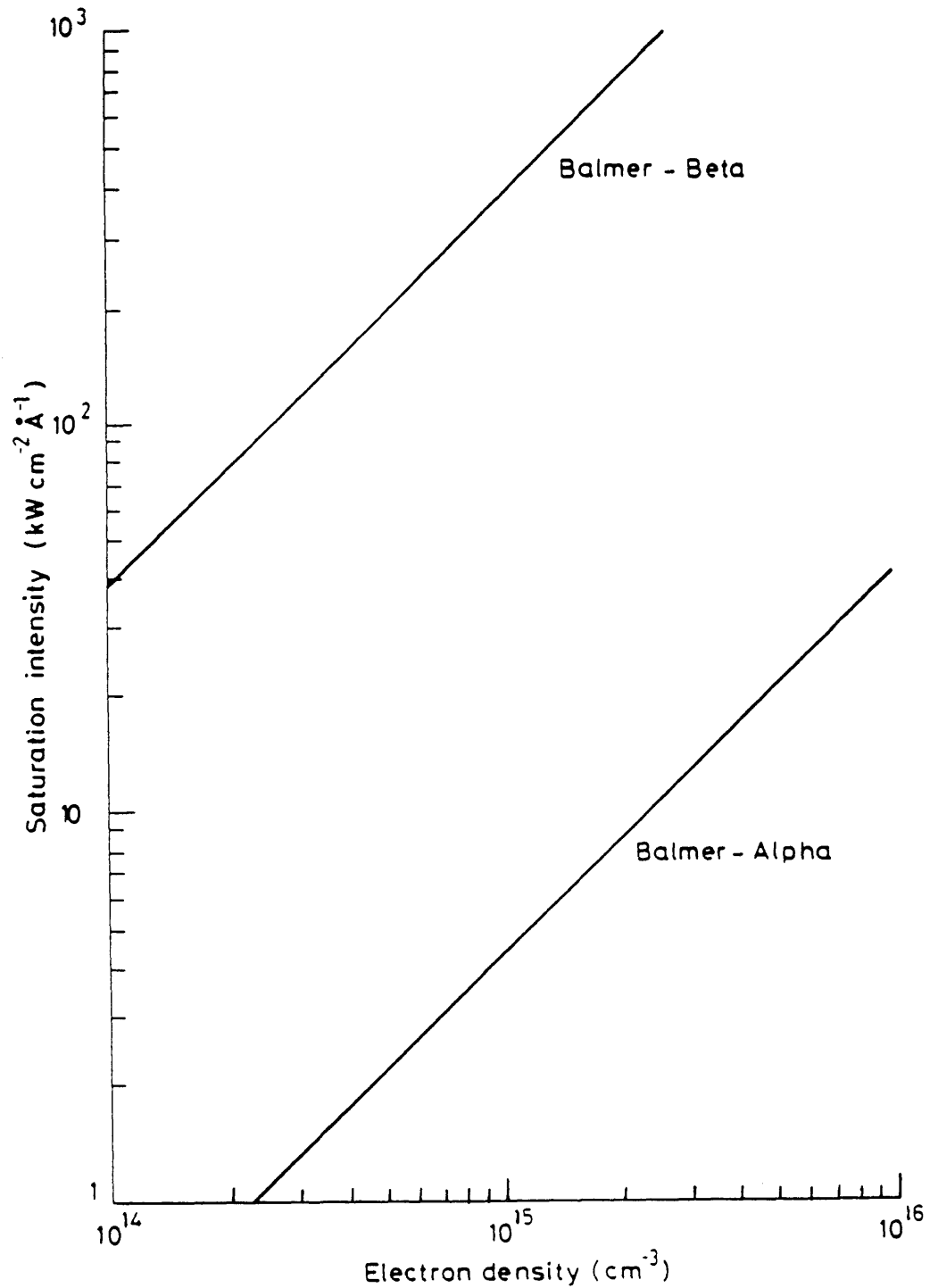


Figure 1.1 The Laser Intensity Required to Saturate the Balmer-Alpha and Beta Transitions as a Function of the Electron Density.

This is for a 50% degree of saturation ($f_s=0.5$). The collision rates of Johnson (1972) are used in the calculation for an electron temperature of 1eV.

over two orders of magnitude), but no laser system has been developed to date that is able to produce such high powers at this wavelength (see Chapter 5 and Gohil and Burgess (1983) for a more detailed discussion).

1.2.3 Modelling the Level Population Behaviour.

1) The Burgess-Skinner Two Level Model. In order to interpret the results, the behaviour of the atomic system must be modelled. A simple two level model from Burgess and Skinner (1974) demonstrates the physics of the system.

Consider a laser tuned to a particular transition between two atomic states j and i (j being the upper level and i the lower). At a time $t=0$, the laser is assumed to completely saturate the transition. Denoting the rate at which the upper level is populated by all states other than the lower level by C_j and the rate at which it is depleted (per upper level atom) by D_j , then at a time $t>0$, the upper level population $n_j(t)$ may be described by:

$$n_j(t) = \frac{C_j + C_i}{(g_i/g_j)D_j + D_i} + \left\{ \frac{n_i + n_j}{1 + (g_i/g_j)} - \frac{(C_i + C_j)}{(g_i/g_j)D_j + D_i} \right\} \exp \left[- \frac{((g_i/g_j)D_i + D_j)t}{1 + (g_i/g_j)} \right] \quad (1-3)$$

where C_i and D_i are the corresponding rates for the lower level. From this model, the fluorescence enhancement observed may have one of two characteristic forms depending on the sign of the *second* term in Equation 1-3. Either the fluorescence signal rises rapidly from the initial background level ($n_j(0)$) and then decays exponentially, leveling out to a quasi-equilibrium plateau level which lasts as long as the laser remains on, or the intensity will simply rise up to an equilibrium value, the characteristic "spike" not being seen. In both cases, the shape of the fluorescence signal will give information on the level populations and the

corresponding electron collision rates. The ratio of the initial spike to the background will essentially give the ratio of the initial pre-laser populations $n_j(0)/n_i(0)$. The time constant of the transient decay will give the combined decay rates of the two levels weighted by their statistical weights. The ratio of the plateau height to the background will give the ratio of the decay rates as measured by the spike decay to $(C_i+C_j)/n_j(0)$. Combination of these results will separately give D_i , D_j , $n_j(0)/n_i(0)$ and $(C_i+C_j)/n_j(0)$.

A typical Balmer-alpha fluorescence trace in hydrogen is shown in Figure 1.2, taken from Burgess et al (1980). The characteristic spike is clearly seen, decaying to a plateau level which remains constant for the duration of the laser pulse.

ii) Collisional Radiative Model. The Burgess-Skiner model may be used to predict the overall behaviour of the two levels under investigation but in the light of significant disagreement between experiment and theory, a more rigorous treatment is required. A 20 level collisional-radiative (CR) code described by Gohil (1982) was used throughout this thesis to analyse results, and is referred to throughout as the CR code. The populations of the first 20 levels in hydrogen are treated explicitly, the code incorporating all the relevant atomic processes.

The equation describing the behaviour of the population of the i^{th} level is:

$$\frac{dn_i}{dt} = \text{Rates into level } i - \text{Rates out of level } i$$

$$\frac{dn_i}{dt} = \sum_{j=i+1}^{\infty} n_j A_{ji} + n_e \left\{ \sum_{\substack{j=1 \\ j \neq i}}^{\infty} n_j C_{ji} + n_e (C_{ci} + C_{ri}) \right\} \quad (1-4)$$

$$- \sum_{j=1}^{i-1} n_i A_{ij} + n_e n_i \left\{ \sum_{\substack{j=1 \\ j \neq i}}^{\infty} C_{ij} + C_{ic} \right\}$$

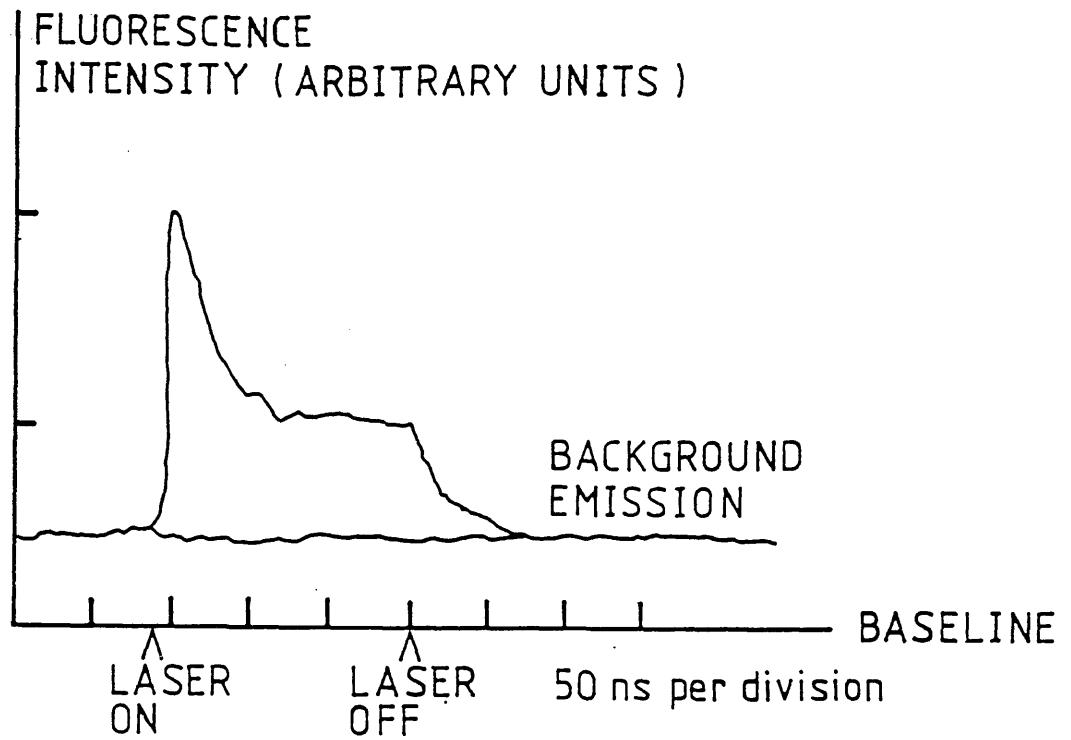


Figure 1.2 Balmer-Alpha Fluorescence in a Hydrogen Plasma.

This shows a typical Balmer-alpha fluorescence trace taken from Burgess et al (1980) for an electron density of $5 \times 10^{14} \text{ cm}^{-3}$ and electron temperature 0.6 eV.

where C_{ic} and C_{ci} are the collisional ionization and recombination rate coefficients to and from the continuum and C_{Ri} is the radiation recombination coefficient into level i and all other symbols have their previously defined and usual meanings. The collision rates are obtained from the semiempirical formulae of Johnson (1972), the spontaneous decay rates from Wiese et al (1966) and the radiative recombination coefficients from Seaton (1959). The electron collision excitation and de-excitation rates are related by the principle of detailed balance,

$$C_{ji} = (g_i/g_j) C_{ij} \exp(E_{ij}/KT) \quad (1-5)$$

The steady state population densities are obtained from Equation 1-4 by setting $dn_i/dt=0$. The perturbation of the level populations as a result of the laser pulse is obtained by incorporating the laser induced rates into the rates matrix and simultaneously solving the rate equations for each level. In this way it is possible to observe the effect of different shaped laser pulses on the level populations. Table 1.1 shows electron collision and radiative rates between the levels of interest in hydrogen for the two extreme plasma conditions studied here. In general, the collision rates (per atom) are larger than the radiative rates.

The quantitative predictions of the behaviour of the level populations from the Burgess-Skinner 2 level model and the 20 level CR code agree to within 20%. The CR code however, gives a more detailed picture since the absolute level populations can be calculated and sensitised (collisionally induced) fluorescence scattering results analysed.

1.3 Introduction to Fluorescence Scattering Measurements in Plasmas.

Laser spectroscopy is a powerful tool for the study of

TABLE 1.1: The Main Electron Collisional Excitation/De-excitation and Radiative Rates Between the Levels of Interest for Visible Fluorescence Work in Hydrogen.

The main rates for the two extreme set of plasma conditions studied are shown. The electron collision rates are calculated from the semiempirical formulae of Johnson (1972) and the radiative rates are taken from Wiese et al (1966). For the 20 us delay time, the plasma conditions are $T_e=1.03$ eV, $n_e=24.6 \cdot 10^{14} \text{ cm}^{-3}$ and for the 70 us delay time, $T_e=0.19$ eV, $n_e=0.38 \cdot 10^{14} \text{ cm}^{-3}$. Note that the radiative rates to the ground state have been reduced by an appropriate escape factor to allow for the fact that the Lyman lines are optically deep in this z-pinch plasma. A_{ji} denotes the radiative rate and K_{ij}/K_{ji} the collisional excitation/de-excitation rate (per atom).

LEVEL		RATE	
n		(s ⁻¹ per atom)	
		20 us	70 us
2	A_{21}	4.7E04	4.7E04
	K_{21}	2.0E07	5.7E05
	K_{23}	2.1E08	2.0E03
	K_{24}	2.3E07	1.4E01
3	A_{31}	5.6E03	5.6E03
	A_{32}	4.4E07	4.4E07
	K_{31}	1.7E06	4.7E04
	K_{32}	6.3E08	1.1E07
	K_{34}	3.1E09	2.3E06
4	A_{41}	1.3E03	1.3E03
	A_{42}	8.4E06	8.4E06
	A_{43}	9.0E06	9.0E06
	K_{42}	7.2E07	1.2E06
	K_{43}	3.3E09	3.6E07
	K_{45}	1.4E10	3.2E07

plasmas, neutral gases and vapours in atomic physics. Early experiments were restricted to the use of fixed frequency lasers. Although not strictly a spectroscopic technique, Thomson scattering (first performed in a plasma by Funfer et al, 1963) has become a routine technique for probing the internal structure of plasmas and is commonly used to determine electron temperatures (see e.g. Ward, 1983). A novel application of this technique was used to measure the pitch angle of the magnetic field lines in a tokamak plasma (Forrest et al, 1978).

However, the development of the tunable dye laser (Sorokin and Lankard, 1966) expanded the applications of laser spectroscopy, and in particular allowed the use of resonant scattering, with its very high cross sections. Many studies of neutral atoms and molecules have been made with this technique (see e.g. Bradley et al, 1970, 1973). Studies in dense atomic vapours have been carried out by Burgess and Eckart (1976) and Salter (1979) in sodium, and by McIlrath and Lucatorto (1977) and Skinner (1980) in barium. An example of how far these techniques have progressed due to the development of dye lasers is given by the work of Pan et al (1980), in which a single sodium atom was detected in a laser beam.

Laser spectroscopy of ionizing and recombining plasmas has only recently been exploited, mainly because the laser photoexcitation rate must compete with relatively rapid electron collision processes. Saturation powers of the order of KWcm^{-2} are required as opposed to mWcm^{-2} for atomic gases and vapours (Burgess, 1978). There is also the problem that there are regions in the ultra violet spectrum which have not yet been reached by laser dyes. Flashlamp excited dyes have been made to work down to $3410 \text{ \AA}$ (Furcumoto and Ceccon, 1970) and certain wavelengths down to $1900 \text{ \AA}$ may be reached by pumping the dye with another laser system (see e.g. 1985 Laser Focus/Electro-Optics Buyers' Guide). Even when a dye is

available at the required wavelength, the laser pulse may not be of sufficiently long duration for the transition to become saturated. Saturation experiments in hydrogen plasmas are thus restricted to transitions with initial level $n=2$ or greater, with the consequent problems associated with interpretation of excited level population behaviour.

1.4 Previous Work in Low Temperature Plasmas.

A variety of work on laser induced fluorescence scattering in plasmas has been performed. In many cases disagreement between experiment and theory has been observed regarding the application of collisional-radiative (CR) models to the behaviour of the perturbed atomic system. There is much evidence that the previously accepted electron-atom collision rates of Johnson (1972) are too high by approximately an order of magnitude for the range of plasma conditions investigated. Work which has attempted to estimate or explain the effect of the collision rates on the system investigated will now be reviewed. This review is a guide to the work presented here on low temperature ($T_e < 1$ eV) laboratory plasmas. A discussion of the fluorescence scattering work on fusion-type plasmas is presented in Chapter 5.

Burrell and Kunze (1972) used the frequency doubled output from a Q-switched ruby laser to pump a dye laser and so observe fluorescence in a low density ($n_e = 3 \times 10^{12} \text{ cm}^{-3}$) helium plasma. The dye laser was tuned to the $2^3P \rightarrow 4^3D$ transition ($\lambda = 4471 \text{ \AA}$). Observation of the induced fluorescence on other collisionally coupled helium lines enabled the authors to estimate the electron collision rates between the excited states, and fairly good agreement was obtained with their theoretical estimates. In a similar plasma, Burrell and Kunze (1978) pumped the $4^1P \rightarrow 4^1S$ and $3^1P \rightarrow 3^1S$ transitions and the total rate out of the upper levels (measured from the decay of the

induced fluorescence scattering) agreed with their theoretical calculations to within experimental error.

Collins et al (1972) pumped the $2^3S \rightarrow 5^3P$ transition ($\lambda=2945 \text{ \AA}$) in the afterglow of a He plasma of low density ($n_e=10^{10} \rightarrow 10^{12} \text{ cm}^{-3}$). The collisional de-population rate of the upper level was determined by observation of the fluorescence transient decay.

Burgess and Skinner (1972) measured the depopulation rate out of the 3^3D level from the decay of the fluorescence enhancement after pumping the $2^3P \rightarrow 3^3D$ transition ($\lambda=5876 \text{ \AA}$). This work was performed on a plasma of higher density than previously investigated ($3 \times 10^{14} \rightarrow 10^{15} \text{ cm}^{-3}$) with a coaxial flashlamp pumped tunable dye laser and showed that the use of short laser pulses was not always necessary to deduce decay rates. Indeed, once saturation had been achieved additional information could be obtained with laser pulse lengths exceeding decay times of interest.

Work has also been performed in hydrogen, and it is this that has led to the work presented in this thesis. Note that the work performed at Imperial College (Skinner (1974), Burgess et al (1978, 1980), Gohil (1982) and Ward (1983)) was performed on the same plasma source: the recombination phase of a z-pinch plasma for which accurate measurements of the relevant plasma parameters were available. The range of electron densities investigated by these authors was $2 \times 10^{14} \rightarrow 1.2 \times 10^{15} \text{ cm}^{-3}$ and the corresponding temperature range was $0.4 \rightarrow 0.8 \text{ eV}$.

Skinner (1974) used a coaxial flashlamp pumped dye laser to saturate the Balmer- α transition ($\lambda=6563 \text{ \AA}$). From the temporal behaviour of the Balmer- α fluorescence, he was able to deduce the decay rate out of the $n=3$ level, which was found to be significantly lower than that predicted by Johnson (1972).

Fluorescence scattering on a lower density ($1.75 \cdot 10^{13} \text{ cm}^{-3}$) and temperature (0.15 eV) hydrogen plasma was performed by Himmel and Pinnekamp (1977). They resonantly pumped the Balmer- α and - β transitions and measured the decay of the fluorescence signal to evaluate the total depopulating rates for the $n=3$ and 4 levels. Simultaneous pumping on the Paschen- γ ($n=3 \rightarrow 6$) and - δ ($n=3 \rightarrow 7$) and the Balmer- α transitions gave the depopulation rates out of the $n=6$ and 7 levels. Calculations from a set of differential equations coupling the atomic levels together and incorporating the rates of Johnson (1972) agreed well with the experimental results. Although the error bars on their data were large (the error in the measured decay rate out of the $n=3$ level being a factor of approximately 5 for example) the functional dependence of data with n is well described by Johnson's rates. This agreement with Johnson for the main depopulating rates out of the $n=3, 4, 6$ and 7 levels is not consistent with work performed at higher densities at Imperial College. However, for the lower densities investigated ($n_e = 0.9 \rightarrow 3.0 \cdot 10^{14} \text{ cm}^{-3}$) the main rates out of the $n=3$ level (as measured by Burgess et al (1980)) have also been found to agree with Johnson to within the experimental error.

However, work by Himmel and Sowa (1983) on a similar plasma showed disagreement with the rates of Johnson. They worked with a low density ($n_e = 2.8 \cdot 10^{13} \text{ cm}^{-3}$, $T_e = 0.16 \text{ eV}$) deuterium plasma. Himmel and Sowa pumped the $n=2 \rightarrow 6$ transition and observed the transient behaviour of the $n=5, 6$ and 7 levels, comparing their observations with calculations based on rate equations with 5 states. Use of the rate coefficients from the semiempirical formulae of Johnson did not explain their data, but use of those from Vriens and Smeets (1980) did compare favourably to within experimental error. Himmel and Sowa also commented on the apparent agreement with Johnson for similar plasma conditions observed by Himmel and Pinnekamp (1977). They pointed out that the laser induced rate into

the $n=6$ level was the same order of magnitude as the spontaneous rate, and that the accuracy of the depopulation rates obtained from Himmel and Pinnekamp was strongly dependent on the accuracy of the initial population densities and the geometric parameters.

Burgess et al (1978) used a coaxial flashlamp pumped dye laser to saturate the Balmer- α transition in the recombination phase of their z-pinch plasma. At the same time, a confocal flashlamp pumped dye laser was tuned to Balmer- β , and with its power well below the saturation intensity of the line, was used to monitor the $n=2$ population. The decay rates out of the $n=2$ and 3 levels were estimated from the fluorescence decay times and the $n=2$ recovery time after laser irradiation. The Balmer- α plateau decay rate was much slower than would be expected using Johnson's values for electron-atom collision rates (except for the lowest density point) being equal to the spontaneous emission rate A_{32} (see also Ward, 1983). At the highest electron density point, the decay rate was approximately a factor of 5 slower than that given by Johnson. The $n=2$ decay rate observed followed the $n=2 \rightarrow 1$ de-excitation rate rather than the dominant $2 \rightarrow 3$ excitation rate. (At an electron density of $1.1 \times 10^{15} \text{ cm}^{-3}$, the $n=2 \rightarrow 1$ rate should be more than a factor of 7 smaller than the $n=2 \rightarrow 3$ rate). Burgess et al also observed saturation at much lower powers than predicted.

Perhaps the most exhaustive work on this topic to date has been done by Burgess et al (1980). Using the same plasma conditions as the previous work, resonant pumping was performed on the Balmer- α transition and observations of the resulting fluorescence on both the Balmer- α and - β transitions were made. The time history of the observed fluorescence traces was compared with the output of a multi-level collisional-radiative (CR) code which incorporated the electron-atom collision rates of Johnson. Large deviations between the experimental

results and theoretical predictions were observed. For example, the observed decay of the Balmer- α fluorescence transient was much slower than that predicted by the code, but agreed well when all the rates (except those associated with the ground state) were reduced by an arbitrary factor of ten. The actual magnitude of the fluorescence enhancement was larger than predicted by factors of between 2-4, and the laser intensity required to saturate the Balmer- β transition was over 50 times lower than theoretically predicted. The steady-state $n=2$ population (i.e. with no laser) was also nearly an order of magnitude higher than theoretically predicted for the higher electron densities studied ($n_e > 8 \times 10^{14} \text{ cm}^{-3}$).

The authors concluded that the anomalies observed were most likely to be due to the rates as calculated from the formulae of Johnson being too large by a significant factor. To account for the enhanced quiescent $n=2$ level population, an extra rate into this level would also be required. The Balmer- α and - β optical depths were too low to account for the discrepancies and the presence of such ions as H_2^+ and H^- were discounted because of the low LTE populations of such ions in the plasma. Burgess et al (1980) also pointed out that no single experimental effect could explain all the discrepancies and that any laser induced effects were orders of magnitude too small to affect the results.

Gohil (1982) observed the risetimes of the Balmer- α and - β fluorescence traces whilst saturating the Balmer- α transition with a fast rising laser pulse (risetime 3.5-4 ns measured by the detection system used for observation of the fluorescence traces). Both the Balmer- α and - β risetimes were larger than predicted by the CR code. The long Balmer- α risetime could be explained by opacity effects, but the Balmer- β result could not. The $n=4$ risetime (with respect to the measured $n=3$ risetime) agreed with theory when all the electron collision rates were reduced by factors of between 5 and 20. The actual

$n=4$ sensitised fluorescence enhancements also agreed well with theory when the depopulating rates only out of $n=4$ were reduced, in this case by a factor of 10.

Work was also performed by Gohil on a different, slightly hotter plasma. Balmer- α scattering was observed at two delay times only, corresponding to conditions of $n_e=0.7 \times 10^{14} \text{ cm}^{-3}$, $T_e=0.7 \text{ eV}$ and $1.1 \times 10^{15} \text{ cm}^{-3}$, 1.5 eV . The fluorescence transient was observed to decay slower than predicted by a factor of approximately 10. The high plateau and peak observed at the higher density and temperature conditions could be explained by a reduction in the collision rates associated with the $n=2$ and 3 levels, but a similar reduction in the rates did not explain the observed enhancement for the lower density studied.

Chapter Two

FLUORESCENCE ANOMALIES IN A LOW TEMPERATURE HYDROGEN PLASMA.

In this chapter, anomalous fluorescence scattering results obtained prior to this work are discussed. Then an alternative (lower) set of electron collisional excitation/ de-excitation rates is incorporated into the CR code which has the effect of eliminating or reducing some of the previously observed discrepancies. Account is also taken of processes involving H^- and H_2^+ ions which populate the $n=2,3$ and 4 atomic levels. However, two clear anomalies still remain unaccounted for: the high $n=3$ and collisionally induced $n=4$ plateau enhancements observed, and the (pre-laser) $n=2$ level density which is still considerably higher than predicted for the earlier time delays in the recombination phase.

2.1 The Discrepancies.

As briefly outlined in Section 1.3, much fluorescence scattering work has been performed at Imperial College on one particular, well diagnosed recombining plasma produced by a z-pinch device run with hydrogen as the filling gas (see e.g. Burgess et al, 1978, 1980). Many discrepancies between theory and experiment have been found as a result of this work and may be summarised as follows:

The theory, as predicted by the 20 level CR code described in Section 1.4 fits:

i) the $n=3, 4$ and 5 peak enhancements (n_i^{PK}/n_i^0) quite well, particularly for the higher temperatures and densities ($T_e > 0.45$ eV, $n_e > 3 \times 10^{14}$ cm $^{-3}$).

ii) the initial pre-laser population ratio of the $n=3$ and 4 level (n_3^0/n_4^0). The fit in this case is not as good as for point (i).

The theory does not fit:

i) the initial pre-laser populations of the $n=2$ level. The experimentally determined populations are in general higher than the predicted values.

ii) the peak H- α fluorescence enhancements for the lower temperatures and densities ($n_e < 3 \times 10^{14}$ cm $^{-3}$, $T_e < 0.45$ eV) which are lower than those predicted from the theory. Ward (1983) showed that the theory was sensitive to the risetime of the laser pulse, however, much of the data was obtained using a fast rising pockel cell switched laser. Ward incorporated a linear risetime of 20 ns for the laser pulse into the CR code but this still produced values that were significantly higher than the experimental data for these low densities.

iii) the plateau ratios n_3^{PL}/n_3^0 , n_4^{PL}/n_4^0 obtained whilst saturating the H- α transition, which are higher than theory.

iv) the n=3 and in particular the collisionally induced n=4 risetimes. The experimentally determined risetimes are slower than predicted.

v) the transient (peak) decay time which is also slower than theory except for the lowest temperatures and densities considered.

vi) the decay of the plateau. This is again slower than theory. In fact the decay rate measured agrees better with the radiative rate A_{32} rather than the main collision rate involved K_{32} .

vii) the measured n=2 decay during H- α irradiation. Only at the lowest temperature and density investigated ($2 \times 10^{14} \text{ cm}^{-3}$, 0.45 eV) does the theory and experiment agree. In all other cases the measured decay rate is lower than theory.

viii) the measured laser saturation power. Both the Balmer-alpha and beta transitions have been observed to saturate at lower powers than predicted for all plasma conditions investigated.

The anomalies mentioned in points (iv), (v), (vi), (vii) and (viii) could be explained if the effective electron collision excitation and de-excitation rates were significantly lower than those predicted theoretically from Johnson (1972) Points (i) and (ii) however, could not be explained only by a decrease in the rates alone. An extra rate into the n=2 and possibly the n=3 level not previously considered in the CR code could explain these discrepancies. To account for the high plateau ratios (point (iii)) a combination between the lower collision rates and the extra rate would be required.

Burgess et al (1980) estimated that the accepted theoretical electron collisional excitation/de-excitation rates of Johnson were too high by approximately an order of magnitude, this being particularly so for the higher temperature and density points investigated. Gohil (1982) pointed out that the extra rate into the $n=2$ level would need to be of the order of $10^{18} \text{ cm}^{-3} \text{ s}^{-1}$ to compare with the dominant electron collision rates populating this level.

The following sections consider i) the effect of an alternative, lower set of electron collision rates on the anomalies and ii) possible mechanisms that could populate the $n=2$ and/or 3 levels.

2.2 Alternative Rate Formalisms.

To calculate the rate coefficients for electron-atom excitation collisions, it is desirable to have an analytical representation of the cross sections that is easy to compute. In general, even for atomic hydrogen, calculations of such cross sections near threshold are quite rare. This is especially the case for transitions between excited levels.

An example of a semiempirical formula based on the most reliable cross section data available at the time is found in Johnson (1972). Johnson produces expressions for cross sections and rate coefficients for electron ionization and collisional excitation (for $n \rightarrow n'$ transitions where $n' > n$) in atomic hydrogen for a wide range of temperatures.

The ionization cross sections from Johnson are in good agreement with results from Born calculations at high energies (Omidvar, 1965) and classical cross sections at low energies (Gryzinski, 1959) and also with the semiempirical results of Sampson and Golden (1970) for

ionization from the ground state.

Johnson's collisional excitation cross sections for transitions from the ground state agree well with other references at low energies (Burke et al 1967, Seaton 1959). For transitions involving the higher n levels, Johnson had very few experimental measurements and theoretical calculations available upon which to base his semiempirical formula, although the impact parameter approximation of Seaton (1962) (which should apply for strong dipole transitions) set upper limits on the $n \rightarrow n+1$ transitions. Indeed, Johnson's $n=2 \rightarrow 3$ collisional cross sections agree well with impact parameter results from the calculations of Saraph (1964). For transitions involving levels with $n > 2$, Johnson estimated that his cross section were in error by less than a factor of 2, and for low and intermediate energies more accurate than the classical and Born cross sections.

Gee et al (1976) used classical theories to produce an alternative set of equations for the calculation of the collisional excitation cross sections for transitions involving $n, n' > 5$ for a range of temperatures $10^6/n^2 < T_e(K) < 3 \cdot 10^9$. This set of cross sections was not incorporated into the CR code however because the electron temperature of the present plasma was outside the range quoted by Gee over which he estimated his formulae to be applicable. Also, the main rates of interest in the present study are for transitions involving $n, n' \leq 5$.

Another set of cross section formulae was produced by Sampson and Golden (1970). For $n \leq 2$ at low energies, their rate coefficients agree well with those from Johnson and reproduce the Born approximation results at high energies. For transitions involving $n > 2$ however, their results for low energies are larger than Johnson. If this set of collisional excitation rates was incorporated into the CR code, even worse agreement than that which exists

already would be obtained.

Callaway and McDowell (1983) in their review of electron collisional cross sections in atomic hydrogen concluded that there was little hope of improving upon the existing results of Johnson for the threshold regime, pointing out the lack of reliable experimental and theoretical atomic data available for the $n \geq 3$ levels.

However, a semiempirical formula given by Vriens and Smeets (1980) based (like Johnson) on previous theoretical calculations and experimental data has been found which produces rate coefficients that are in general, approximately a factor of 3 smaller than those given by Johnson near threshold. The electron ionization rates as given by Vriens and Smeets are also slightly different to Johnson for the excited levels. For example, at a temperature of 1 eV, the $n=2$ ionization rate coefficient given by Vriens and Smeets is a factor of 1.6 larger than that given by Johnson, this factor decreasing with increasing n until for $n > 16$, Johnson's ionization rate coefficients become slightly larger.

The electron collisional excitation coefficients for the $n=1 \rightarrow 2$ transition are typically a factor of 9 lower than those given by Johnson and for the $n=1 \rightarrow 3$ transition are typically a factor of 4 lower for $T_e < 2.15$ eV. For transitions between excited levels the rate coefficients are approximately a factor of 3 lower for the threshold regime. Tables 2.1-2.3 show a comparison between the two sets of rate coefficients (cm^3s^{-1}) for the transitions of interest and the range of temperatures in the 70 cm z-pinch plasma. In all cases the coefficients (and therefore the actual rates per atom) from the semi-empirical formulae of Vriens and Smeets are lower than those of Johnson.

Figure 2.1 shows a flow diagram of the modifications made to the CR code by the author. Previous studies did not

TABLES 2.1-2.3: Rate Coefficients Between the Main Transitions of Interest in the 70 cm Z-Pinch.

These tables show the main rate coefficients associated with the fluorescence work from the semiempirical formulae of 1) Johnson (1972) and 11) Vriens and Smeets (1980). The ratio of Johnson's to Vriens' rate coefficients is also shown. Note that the symbol E-x denotes 10^{-x} .

TABLE 2.1: The n=2-3 Rate Coefficient (C_{23}).

Temperature T_e (eV)	Rate Coefficient (cm^3s^{-1})		
	Johnson	Vriens	Ratio
0.2	5.2E-11	1.2E-11	4.4
0.5	1.3E-8	3.6E-9	3.5
1.0	8.7E-8	2.9E-8	3.1
1.5	1.7E-7	6.1E-8	2.9

TABLE 2.2: The n=3-4 Rate Coefficient (C_{34}).

Temperature T_e (eV)	Rate Coefficient (cm^3s^{-1})		
	Johnson	Vriens	Ratio
0.2	6.2E-8	2.0E-8	3.1
0.5	5.1E-7	2.0E-7	2.6
1.0	1.2E-6	5.0E-7	2.5
1.5	1.8E-6	7.5E-7	2.4

TABLE 2.3: The n=4-5 Rate Coefficient (C_{45}).

Temperature T_e (eV)	Rate Coefficient (cm^3s^{-1})		
	Johnson	Vriens	Ratio
0.2	8.5E-7	3.3E-7	2.6
0.5	3.0E-6	1.3E-6	2.4
1.0	5.6E-6	2.4E-6	2.3
1.5	7.4E-6	3.2E-6	2.3

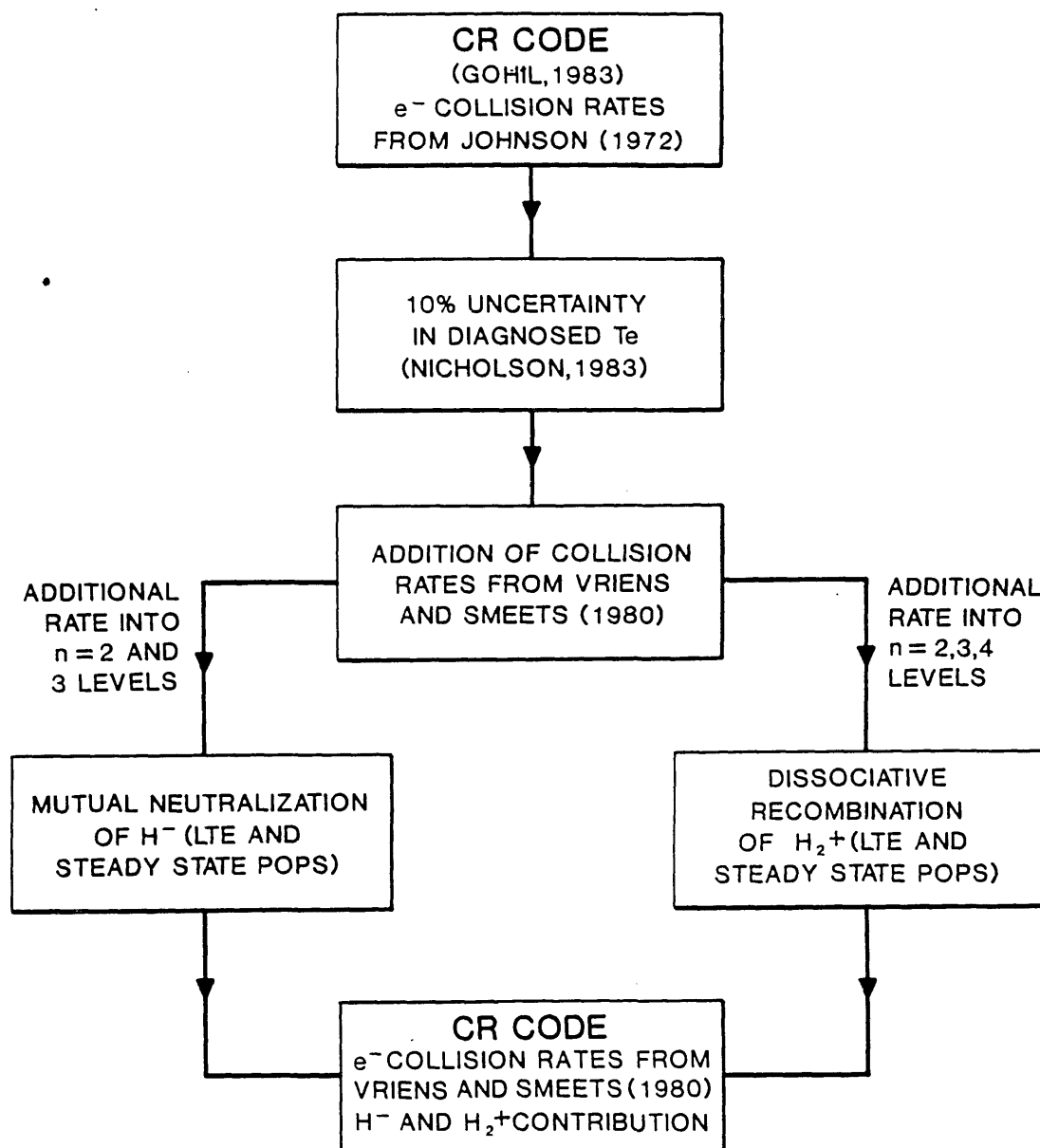


Figure 2.1 Flow Diagram of Modifications Made to the CR Code.

This flow diagram shows the modifications that have been incorporated into the CR code by the author.

include any estimate of the error associated with the theory (as predicted from the CR code). In the following sections, the error in the theory has been calculated from a 10% uncertainty in the diagnosed electron temperature (Nicholson, 1983). The errors associated with the electron density diagnostic are much smaller being approximately 3% (Nicholson, 1983). Then an alternative (lower) set of electron collision rates (due to Vriens and Smeets) was incorporated into the code. Finally, the populations of H_2^+ and H^- ions were calculated and the extra rates into the $n=2,3$ and 4 levels (due to dissociative recombination and mutual neutralisation from these ions) were also incorporated into the CR code. The output of the code was then compared to experimental results (obtained by BMSW and other workers) to see the effect of these changes. This is discussed in the sections below.

2.2.1 Comparison Between Alternative Rate Coefficients and Previous Work. Previous experimental results from Skinner (1974), Burgess et al (1978, 1980), Gohil (1982) and Ward (1983) were compared with results from the CR code using the rates of Vriens and Smeets (1980). This was to determine whether the output thus obtained was a better description of the data than that for the rates of Johnson (1972). No observable difference to the temporal behaviour of the level populations was observed when the ionization rates of Vriens and Smeets were used instead of those of Johnson (whilst using the collision rates of Johnson). Any difference in the output from the CR code was thus attributed solely to the different excitation/de-excitation rates.

1) The $n=2$ Level Population. Figure 2.2 shows the measured $n=2$ population in the absence of the laser pulse, n_2^0 from Burgess et al (1980) and the predicted $n=2$ density from LTE (Saha Boltzmann) calculations and from the CR code (steady state) using the rates of Johnson and Vriens and Smeets. The CR code output

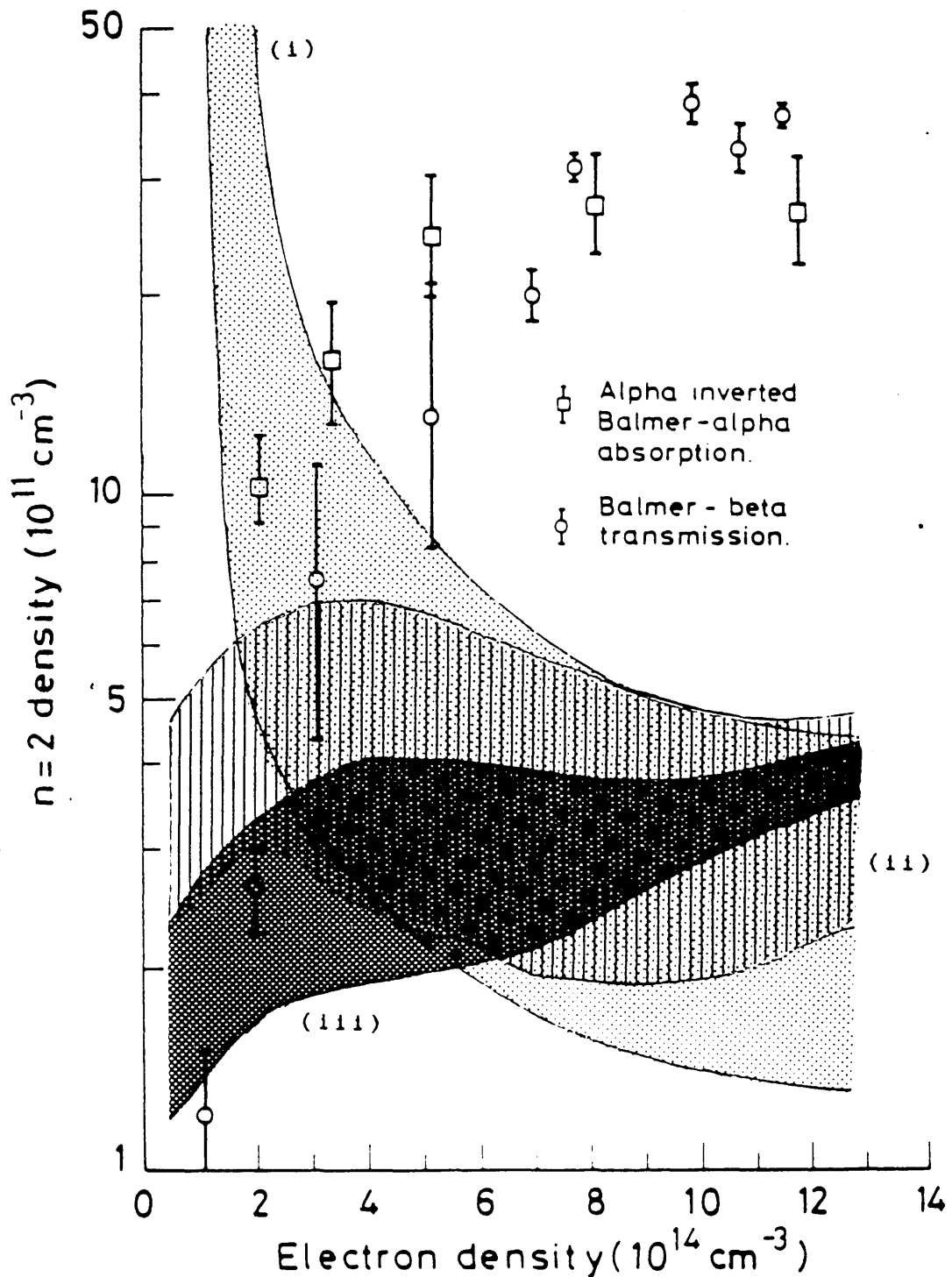


Figure 2.2 The Steady State $n=2$ Level Population as a Function of Electron Density.

The experimental points are reproduced from Burgess et al (1980) and the continuous regions are the theoretical $n=2$ densities as predicted from the CR code for:

- i) LTE (Saha Boltzmann)
- ii) the electron collision rates of Johnson (1972)
- iii) the electron collision rates of Vriens and Smeets (1980).

overlaps the LTE region at the higher temperatures and densities showing the dominance of the electron collision rates over the radiative rates in this region. In general, the measured $n=2$ density is much higher than that predicted by theory for $n_e > 3 \times 10^{14} \text{ cm}^{-3}$, the discrepancy increasing with temperature and density. For $n_e \geq 8 \times 10^{14} \text{ cm}^{-3}$, the measured $n=2$ density is approximately an order of magnitude higher than theory. The predicted density, as calculated from the rates of Vriens and Smeets is not much different to that calculated from the rates of Johnson, showing that a stable equilibrium between the level populations exists, being fairly independent of the exact collision rates chosen.

Gohil (1982) showed that a reduction in the excitation rates K_{23} , K_{34} and K_{45} (but not the corresponding de-excitation rates) by an arbitrary factor of 10 gave a better approximation to the measured $n=2$ density. However, this would require a breakdown of the detailed balance assumption and there is no evidence to suggest that this is so.

ii) The Peak Enhancement. Figure 2.3 shows the H- α peak enhancement (n_3^{PK}/n_3^0) as measured by Burgess et al (1980) as a function of electron density. Also shown are the results from the CR code using i) the collision rates of Johnson and ii) the rates of Vriens and Smeets. The overall distribution of the data points is not described well by the output from the CR code for either set of collision rates. For the higher density points ($n_e > 3 \times 10^{14} \text{ cm}^{-3}$) the results from Johnson are a fair approximation to the data and the results from Vriens and Smeets are too high. Neither set of results approximates well to the lower temperature and density points ($n_e < 2 \times 10^{14} \text{ cm}^{-3}$, $T_e < 0.38 \text{ eV}$) even allowing for a $\pm 15\%$ systematic error in the data reduction estimated by Ward (1983).

The peak height essentially represents the ratio of the initial pre-laser level populations n_2^0/n_3^0 as described

in Section 1.2. As discussed in the previous subsection, the $n=2$ density is known to be much higher than predicted by theory for these higher densities, and so the absolute $n=3$ density must also be higher than theory for these conditions (since the level population ratio is described well by Johnson).

Gohil (1982) measured the collisionally induced H- β peak enhancement in a pure hydrogen plasma and also in a hydrogen in helium plasma (5% hydrogen by volume) whilst saturating the H- α transition. The data is shown in Figure 2.4. For the pure hydrogen case, the predicted values from the CR code for i) the rates of Vriens and Smeets and ii) the rates of Johnson reduced by a factor of 10 (except for those rates associated with the ground state) are also shown. Output from the CR code for the unaltered rates of Johnson is not shown in this case because it agrees with the output for the Vriens and Smeets rates. The theory from the rates of Vriens and Smeets (and indeed Johnson) gives good agreement with the data. The theory from the reduced rates of Johnson gives no agreement with the data at all, the measured peak enhancements being much higher than predicted.

For the hydrogen in helium case also shown in Figure 2.4, the predicted values of the peak enhancement are shown from the CR code for the rates of i) Johnson and ii) Vriens and Smeets. For his analysis, Gohil assumed the temperature of this plasma to be approximately 0.1 eV higher in this range of electron densities than for the pure hydrogen case, based on previous temperature measurements at higher densities (Mahon, 1973). He concluded that the theory as predicted by the CR code for Johnson's rates fitted the data points well. However, Nicholson (1983) diagnosed the temperature to be 0.26 eV hotter than the pure hydrogen case in this region. Using the temperature as diagnosed by Nicholson the results from the CR code using the rates of Johnson are lower than the experimental data, and the enhancements

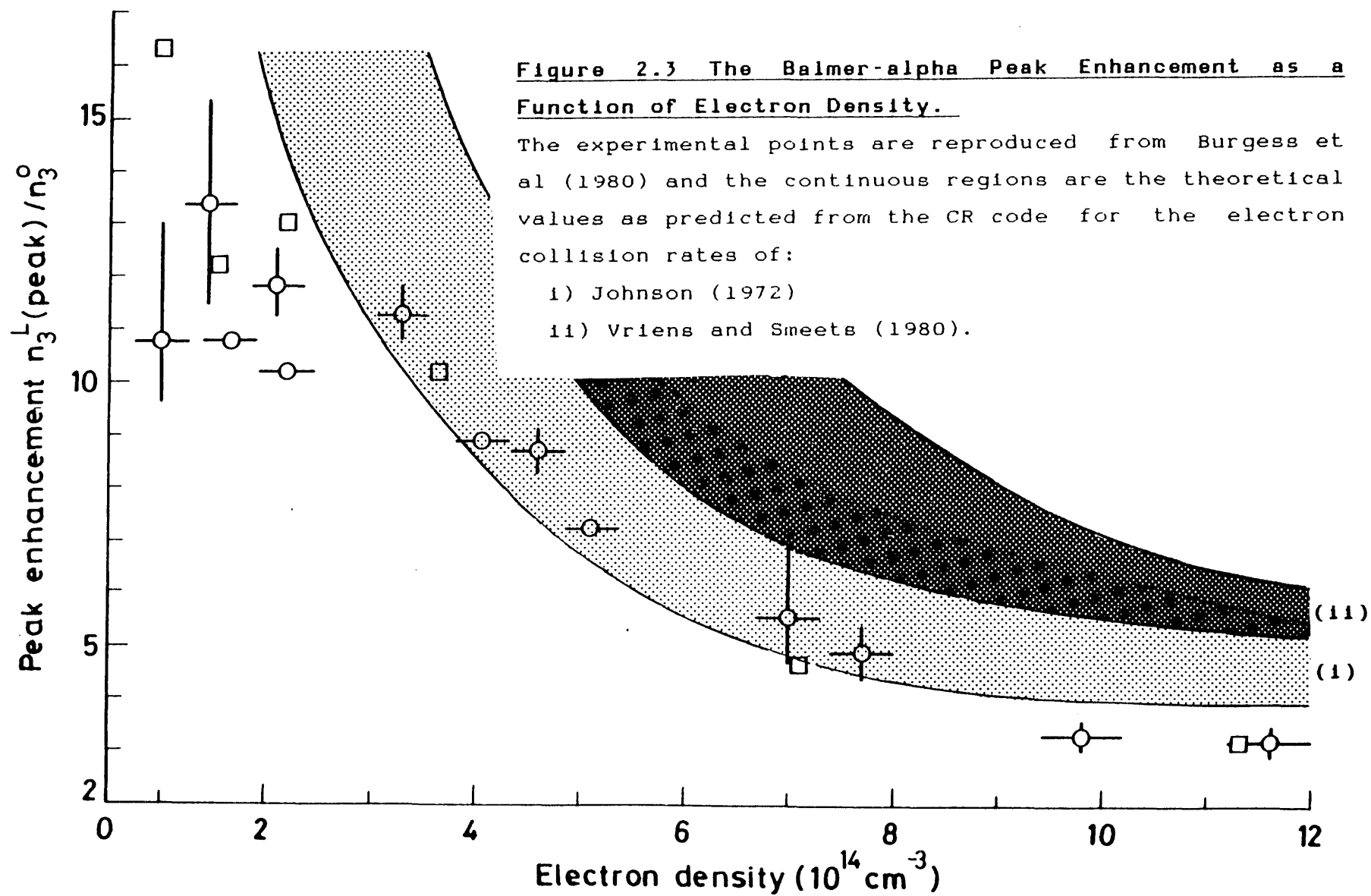
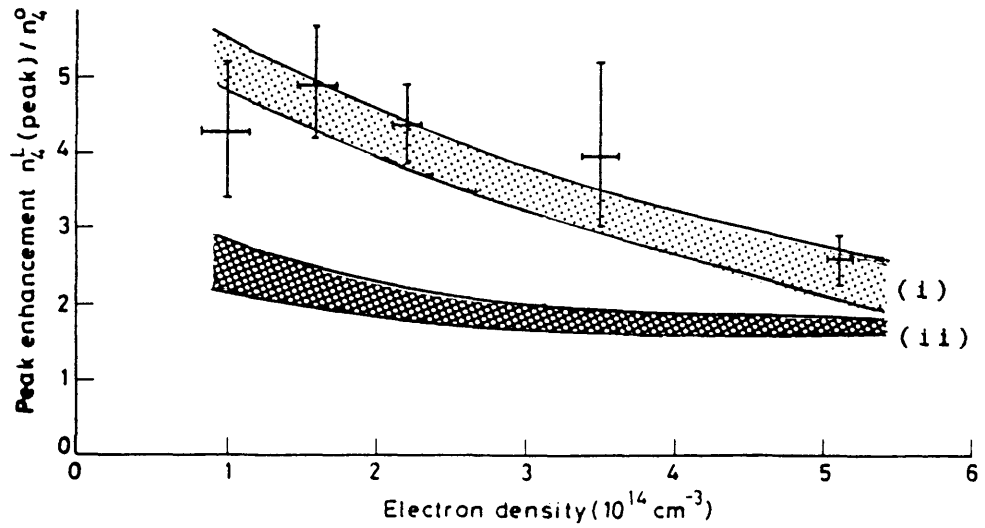


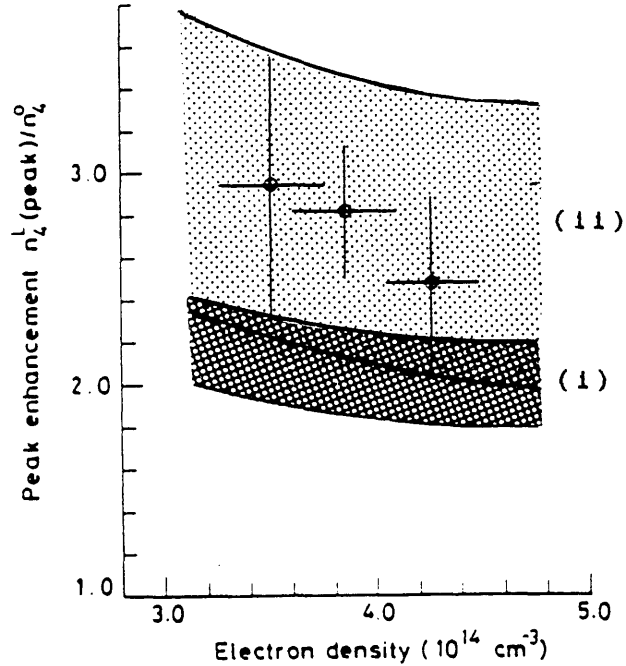
Figure 2.4 The Collisionally Induced Balmer-beta Fluorescence Peak Enhancement as a Function of Electron Density.



a) Pure Hydrogen Plasma.

The experimental points are reproduced from Gohil (1982) and the continuous regions are the theoretical values as predicted from the CR code for the collision rates of:

- 1) Vriens and Smeets (1980)
- 11) Johnson (1972) reduced by a factor of 10.



b) Hydrogen in Helium Plasma (5% by Volume).

The experimental points are as above and the theoretical values as predicted from the CR code for the collision rates of:

- 1) Johnson (1972)
- 11) Vriens and Smeets (1980).

predicted from the rates of Vriens and Smeets are a good fit to the data. The theory is more sensitive to the electron collision rates than in the case of the pure hydrogen plasma because of the larger rate coefficients (in particular the C_{34}) at these higher temperatures.

The collisionally induced H- β peak enhancement is in part sensitive to the H- α peak enhancement but is more sensitive to the collisional coupling to the laser dependent n=3 level and the competing processes upwards to the higher levels. Thus a reduction in the rates used in the CR code results in an increase in the predicted H- α peak enhancement but a decrease in the predicted H- β peak enhancement.

iii) The Plateau Enhancement. Figure 2.5 shows the H- α quasi-equilibrium plateau enhancement as measured by Burgess et al (1980) as a function of electron density. Also shown are the predicted values from the CR code using the rates of i) Vriens and Smeets and ii) Johnson reduced by a factor of 10. The predicted enhancements from the unaltered rates of Johnson are not shown (to preserve clarity), the results from this set of rates giving values that differ by approximately 20% to those given by the rates of Vriens and Smeets. A first glance at Figure 2.5 shows that neither set of collision rates produces results which approximate well to the experimental results across the entire range of plasma conditions investigated. For the higher temperatures and densities ($n_e > 2 \cdot 10^{14} \text{ cm}^{-3}$, $T_e > 0.38 \text{ eV}$) the reduced rates of Johnson are a better approximation to the data. For densities below $2 \cdot 10^{14} \text{ cm}^{-3}$, the predicted values from consideration of the Vriens and Smeets' (and indeed Johnson's) rates are a good approximation, the reduced set of Johnson rates giving values which are too large when compared to the experimental data.

Gohil (1982) also observed the H- β plateau enhancement whilst saturating the H- α transition. His data is

reproduced in Figure 2.6 for a pure hydrogen plasma and a hydrogen in helium (5% hydrogen by volume) plasma. In the case of the pure hydrogen case, the predicted values from the CR code are also shown, using the rates of i) Vriens and Smeets and ii) Johnson reduced by a factor of 10. Output from the code is not shown from the unaltered rates of Johnson because there is much overlap with the results from the Vriens and Smeets rates. In general, the experimental data lies above the predicted values for both set of collision rates, with the exception of the data point at $n_e = 1 \times 10^{14} \text{ cm}^{-3}$ which just overlaps the output from the reduced Johnson rates. The data at densities 2.2 and $3.5 \times 10^{14} \text{ cm}^{-3}$ disagree most with theory, being higher than the predictions of the CR code from the reduced rates by factors of approximately 1.5 and 1.8.

Figure 2.6 also shows Gohil's data for a hydrogen in helium plasma. In this case, the predicted values from the CR code are shown using the rates of i) Johnson ii) Vriens and Smeets and iii) Johnson reduced by a factor of 10. As in the pure hydrogen case, the experimental data lies above the predicted values. The results from the reduced rates of Johnson lie closest to the data. In this case, the data points are approximately factors of between 1.4 and 1.5 larger than theory.

iv) The Fluorescence Risetimes. Gohil (1982) measured the 10-90% risetime of the $n=4$ ($H-\beta$) sensitised fluorescence signal whilst saturating the $H-\alpha$ transition. He measured the laser risetime to be 3.5-4 ns, the $H-\alpha$ risetime to be 5.6 ± 0.8 ns and the corresponding $H-\beta$ risetime to be in the range 15.3 ± 2.0 to 8.0 ± 0.9 ns for the plasma conditions investigated. In his analysis, he found that the measured risetimes of the $n=4$ level populations were too large by factors of between 1.4 and 1.8 when compared with the output of the CR code (for the rates of Johnson) assuming an absolute $n=3$ risetime of 5.8 ns.

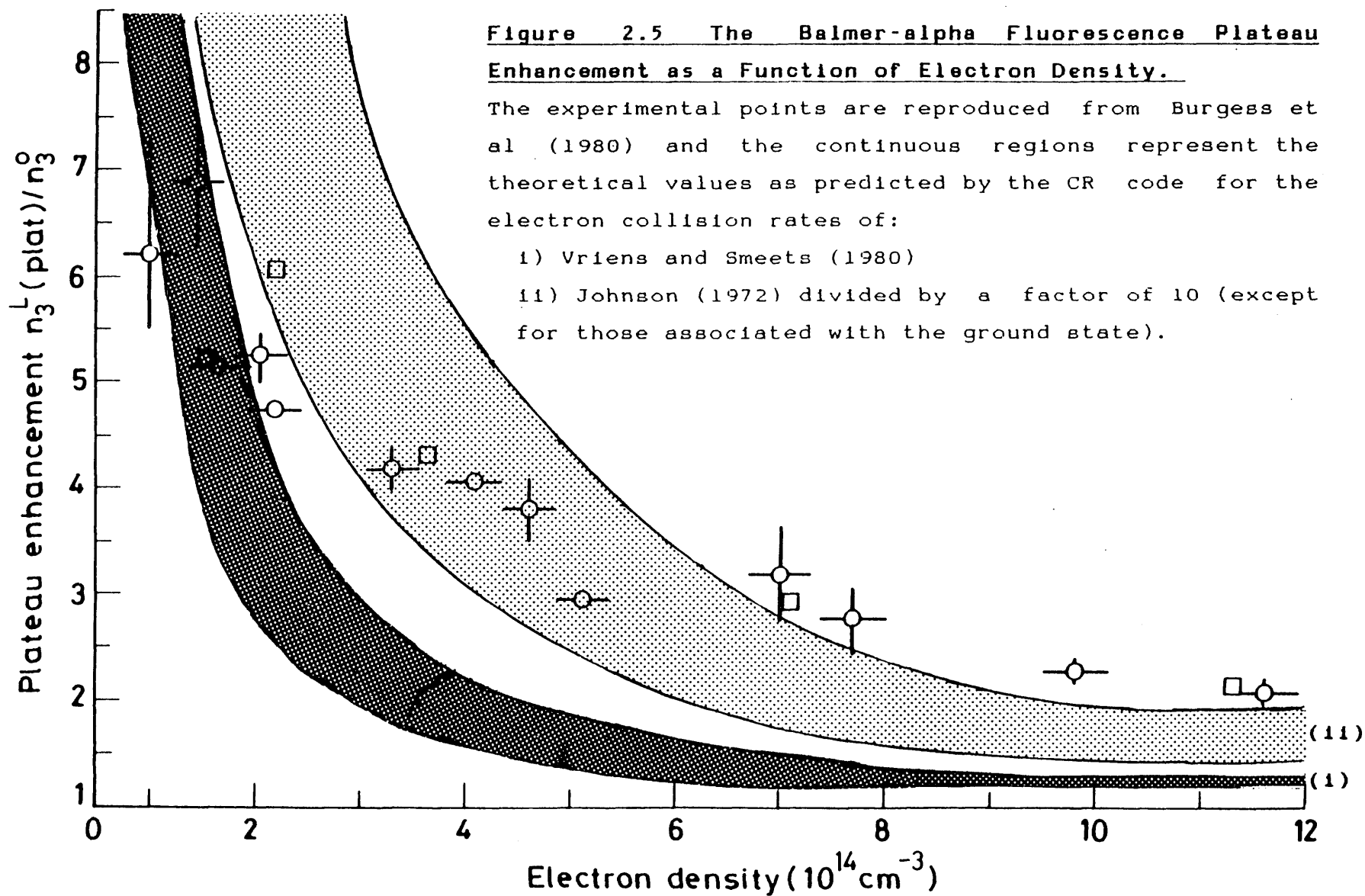
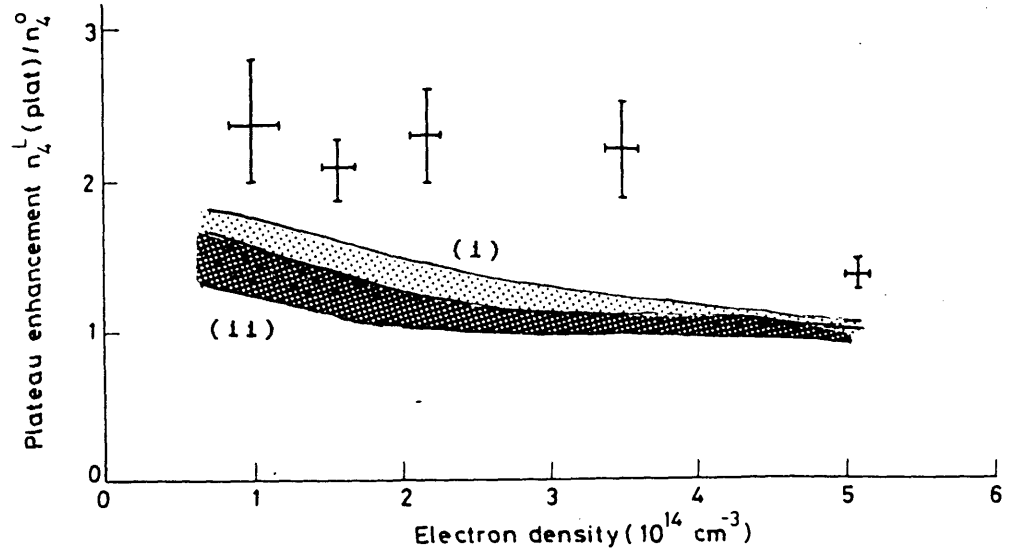


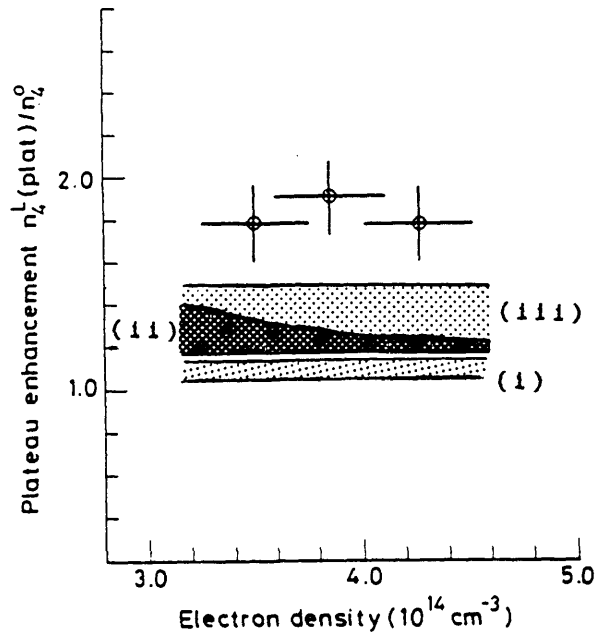
Figure 2.6 The Collisionally Induced Balmer-beta Fluorescence Plateau Enhancement as a Function of Electron Density.



a) Pure Hydrogen Plasma.

The experimental points are reproduced from Gohil (1982) and the continuous regions represent the theoretical values as predicted from the CR code for the electron collision rates of:

- 1) Vriens and Smeets (1980)
- 11) Johnson (1972) reduced by a factor of 10.



b) Hydrogen in Helium Plasma (5% by Volume).

The experimental points are as above and the theoretical values as predicted from the CR code for the electron collision rates of:

- 1) Johnson (1972)
- 11) Vrien and Smeets (1980)
- 111) Johnson reduced by a factor of 10.

However, Gohil took no account of the impulse response function (instrument function) of the experimental system. This was not insignificant since the fluorescence signals had to traverse 30m of RG58 coaxial cable before being digitised. He had noted that the system degraded the 10-90% risetime of a pulse generator (with an actual risetime < 1 ns) to 3.2 ns. This was assumed to be the risetime of the instrument function (in the absence of any other information) so that the data could be analysed.

The instrument function was assumed to be Gaussian in shape so that simple approximations could be used in the convolution of the instrument function with the absolute fluorescence risetime. When this was convolved with an $n=3$ Gaussian risetime of 4.1 ± 0.9 ns the measured $n=3$ risetime of 5.6 ± 0.8 ns was obtained. Assuming the figure of 4.1 ± 0.9 ns to be the absolute $n=3$ risetime, the corresponding $n=4$ risetime from the CR code was then convolved with the instrument function to give the theoretically observed $n=4$ risetime.

Figure 2.7 shows the original data from Gohil, plus the recalculated theoretical risetimes using the rates of 1) Johnson and 11) Vriens and Smeets. The predicted risetimes from the rates of Johnson are lower than the data particularly for the higher temperature and density points. The three data points for the lower temperatures and densities ($1 \cdot 10^{14} \leq n_e \leq 2.2 \cdot 10^{14} \text{ cm}^{-3}$, $0.29 \leq T_e \leq 0.39$ eV) are described well by the CR code output using the rates of Vriens and Smeets. The data points at $n_e = 3.5$ and $5.1 \cdot 10^{14} \text{ cm}^{-3}$ are a factor of approximately 1.4 larger than the theoretical risetimes.

The H- β risetimes are dominated by the $n=3 \rightarrow 4$ electron collision rate K_{34} . The data presented in Figure 2.7 could be taken as direct evidence that K_{34} is nearer to the lower rate as given by Vriens and Smeets than that given by Johnson, at least for the lower temperature and

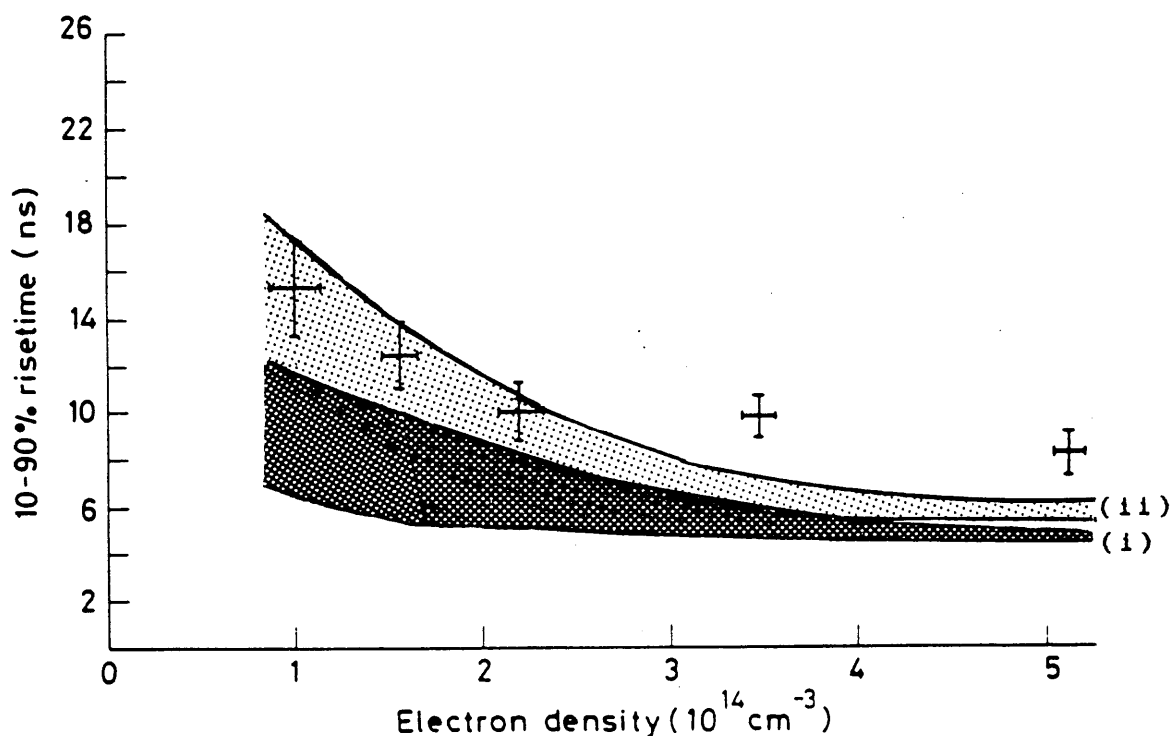


Figure 2.7 The 10-90% Risetime of the Collisionally Induced Balmer-beta Fluorescence as a Function of Electron Density.

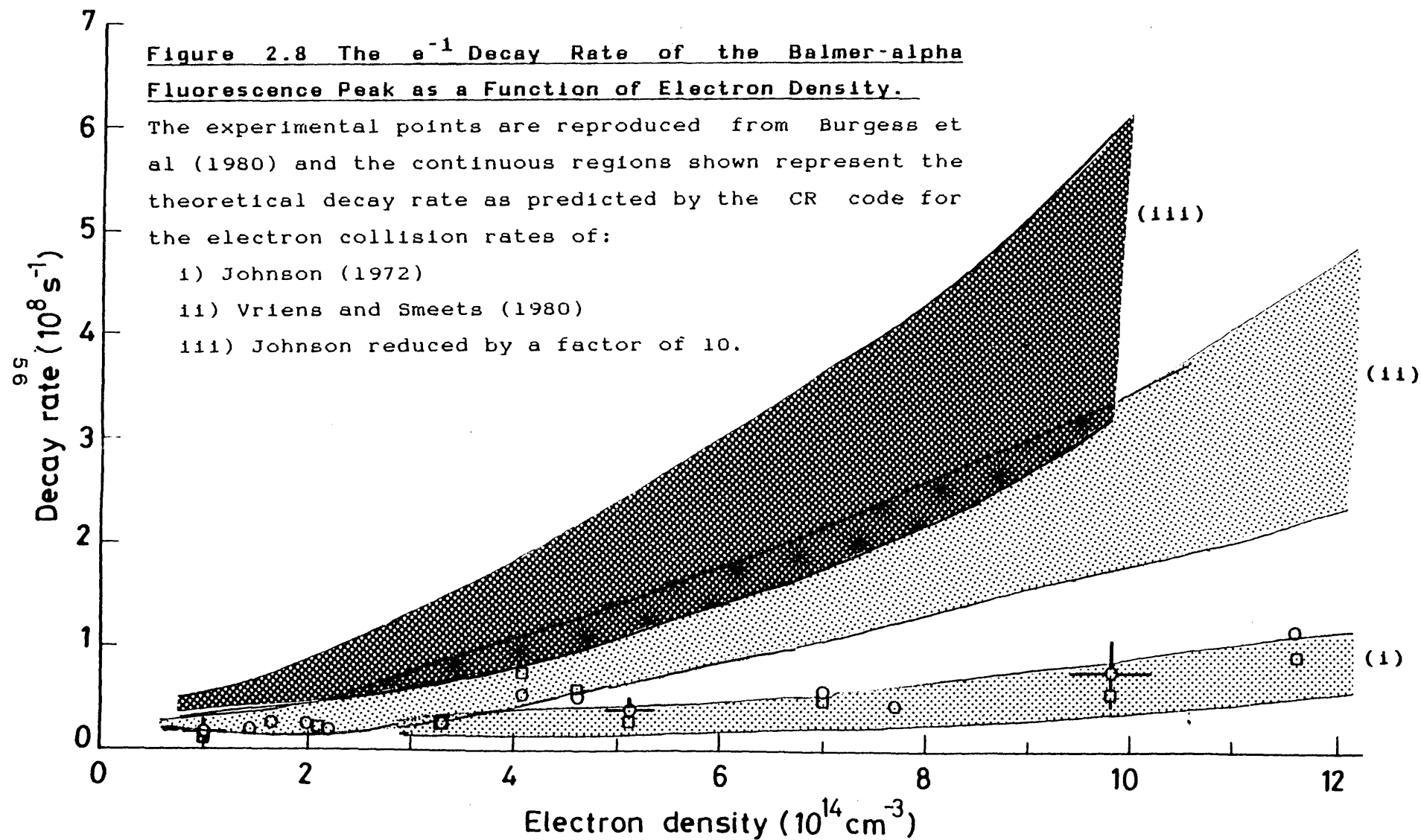
The experimental points are reproduced from Gohil (1982) and the predicted risetimes calculated from the convolution between the 3.2 ns instrument function risetime and the output from the CR code for the electron collision rates of:

- 1) Johnson (1972)
- 1) Vriens and Smeets (1980).

density range of plasma conditions. The two data points at $n_e = 3.5$ and $5.1 \cdot 10^{14} \text{ cm}^{-3}$ are even larger than the predicted risetimes when Johnson's rates are reduced by a factor of 10, suggesting that the $n=3 \rightarrow 4$ rate may be more than a factor of 10 smaller for these plasma conditions.

However, any conclusions drawn from this data are tentative ones only, the analysis as presented here being only an approximation to account for the degradation of the risetimes due to the experimental setup. No error has been assumed in the instrument function risetime and the shape of this pulse may not be Gaussian as assumed. In particular, the 50-90% rising edge of the pulse may be artificially lengthened by the Skin effect (see Dreher, 1967) so that the definition of the risetime as being the 10-90% risetime in the analysis may not be ideal. The agreement shown with the CR code output for the Vriens and Smeets rates is consistent with the measured $n=4$ peak enhancement however, which is also sensitive to the $n=3 \rightarrow 4$ collision rate (see Section 2.2.1 ii)).

v) The Transient (Peak) Decay. Figure 2.8 shows data from Burgess et al (1980) for the e^{-1} decay of the initial H- α transient (the characteristic fluorescence spike) to the plateau level, which is essentially governed by the collision rates out of the $n=3$ level to the higher levels. The predicted decay rate is also shown from the CR code for the rates of i) Johnson ii) Vriens and Smeets and iii) Johnson reduced by a factor of 10. The actual decay rate was observed to increase with temperature but at a much slower rate than predicted by either Johnson or Vriens and Smeets. Output from the CR code for the Vriens and Smeets rates is closer to the experimental data than for the Johnson rates, although for densities larger than $4.6 \cdot 10^{14} \text{ cm}^{-3}$ the measured decay rate is too low even for this set of rates and the output employing the reduced Johnson rates is a better description of the data in this case.



vi) Decay of Plateau. If the laser pulse is terminated sharply in time, the plateau enhancement decay will be dominated by the main rates out of the upper level, these being the $n=3 \rightarrow 2$ rate K_{32} and the radiative rate A_{32} ($K_{32} > A_{32}$ for all but the lowest temperatures and densities investigated). Figure 2.9 shows the measured decay from Burgess et al (1978) plus the theoretically predicted decay from the CR code using the rates of i) Johnson and ii) Vriens and Smeets. The experimental decay rate was observed to be equal to the spontaneous decay rate A_{32} . The results calculated from the rates of Vriens and Smeets are closer to the data points than those calculated from Johnson, but still give decay rates that are too high (by factors of between 1.4-3.0) for electron densities above $5 \times 10^{14} \text{ cm}^{-3}$. Consider the data point at $n_e = 11 \times 10^{14} \text{ cm}^{-3}$ (where the maximum discrepancy occurs). The C_{32} rate coefficient must be equal to or less than $5.46 \times 10^{-8} \text{ cm}^3 \text{ s}^{-1}$ to account for the observed decay rate. However, the value calculated employing Vriens and Smeets's rates in the CR code is 7.8×10^{-8} , and that using Johnson is $2.52 \times 10^{-7} \text{ cm}^3 \text{ s}^{-1}$.

vii) Decay of the $n=2$ Level. Burgess et al (1978) measured the e^{-1} decay rate of the $n=2$ level whilst saturating the $H-\alpha$ transition. Figure 2.10 shows this data and the theoretical decay rate (the main collision rate out of the $n=2$ level, the $n=2 \rightarrow 1$ plus the $n=2 \rightarrow 3$ rate) from the CR code using the rates calculated from the semi-empirical formulae of i) Johnson and ii) Vriens and Smeets. No horizontal error bars were quoted by Burgess et al, but an error of $\pm 1 \text{ us}$ in determination of the exact delay time gives an error of $\pm 8.3\%$ in the electron density determination. Figure 2.10 shows that for the lower density points, ($n_e \leq 5 \times 10^{14} \text{ cm}^{-3}$) the data lies within the range of decay rates predicted the set of rates given by Johnson. For the higher temperatures and densities the data lies within the range of decay rates given by Vriens and Smeets, with the exception of the point at $n_e = 15.8 \times 10^{15} \text{ cm}^{-3}$, which is just below this

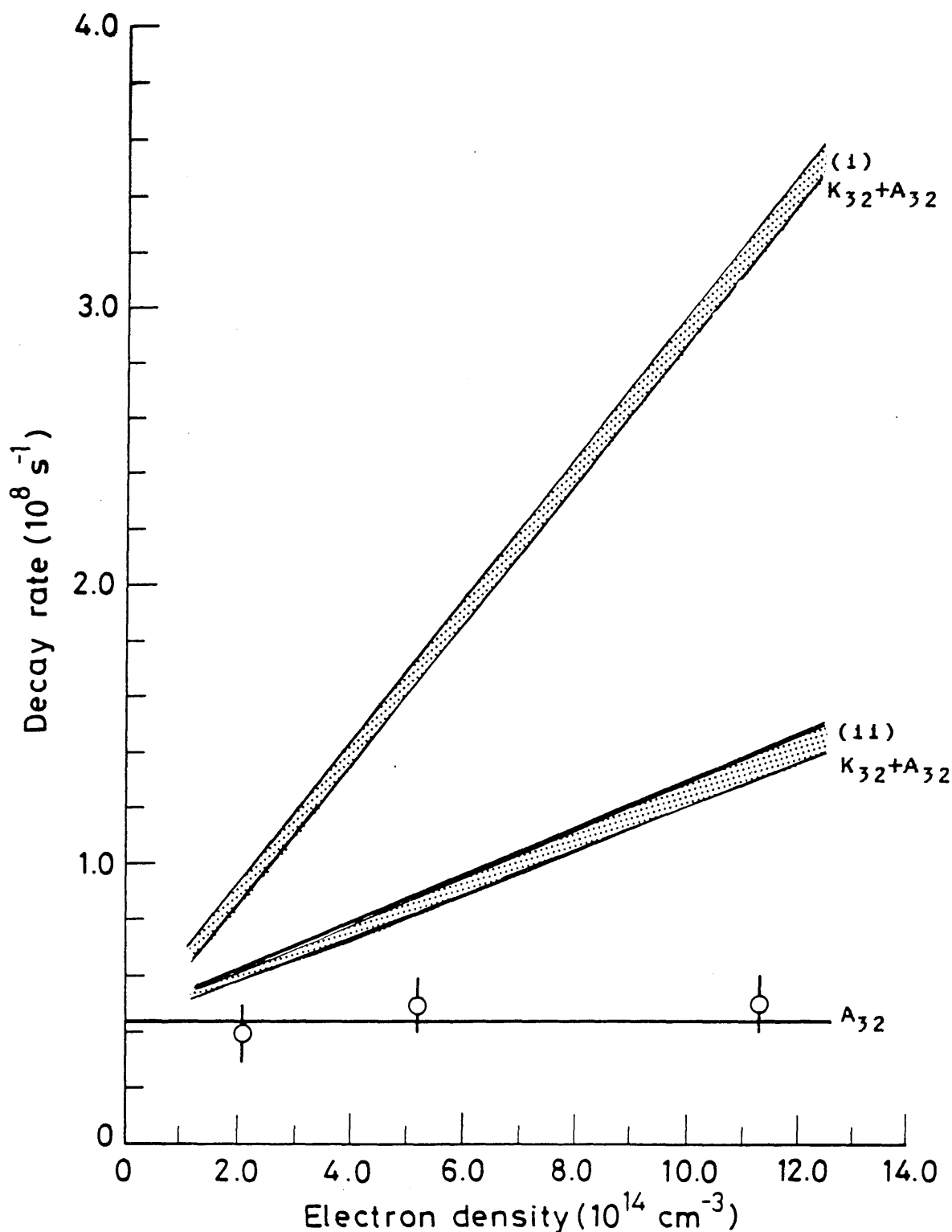


Figure 2.9 The e^{-1} Decay Rate of the Balmer-alpha Fluorescence Plateau as a Function of the Electron Density.

The experimental points are reproduced for Burgess et al (1980) and the predicted values are taken from the CR code for the electron collision rates of:

- i) Johnson (1972)
- ii) Vriens and Smeets.

Also shown is the $n=3-2$ spontaneous decay rate (A_{32}).

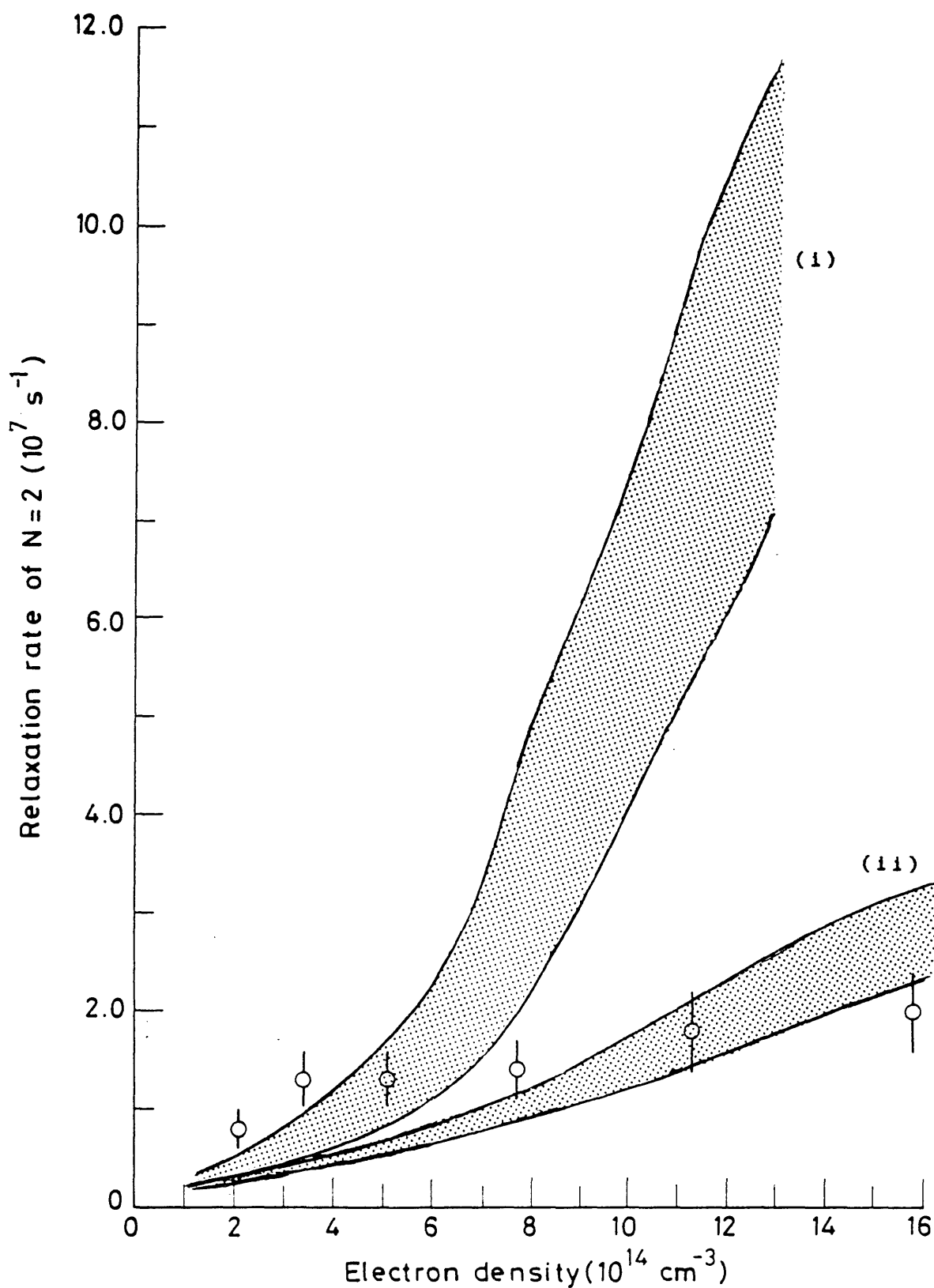


Figure 2.10 The Relaxation Rate of the $n=2$ Level Whilst Saturating the Balmer-alpha Transition as a Function of Electron Density.

The experimental points are reproduced from Burgess et al (1978) and the continuous regions are the theoretical rates from the CR code for the electron collision rates of:

1) Johnson (1972)

11) Vriens and Smeets (1980).

range. This suggests that the true rate out of the $n=2$ level at this point is lower than that given by Vriens and Smeets. Addition of the $n=2 \rightarrow 4$ rate into the calculation does not affect this conclusion since this rate is much smaller than the $n=2 \rightarrow 3$ and $2 \rightarrow 1$ rate (by at least an order of magnitude). Addition of this rate into the calculation makes less than 5% difference at the highest temperature point to the predicted decay rate.

2.2.2 Non-Maxwellian Electron Velocity Distribution.

Gohil (1982) suggested that a non-Maxwellian electron velocity distribution could result in a reduction in the excitation rates relative to the de-excitation rates and that this could account for the high $n=2$ level population observed and also the high H- α fluorescence plateau enhancements observed for $n_e > 3 \times 10^{14} \text{ cm}^{-3}$. However, this hypothesis would produce new disagreements with respect to the $n=4$ and 3 peak enhancements (which agree with output from the code employing the rates of Vriens and Smeets, as discussed in Section 2.2.1 (iii)), and also worsen the agreement between theory and experiment for the decay of the $n=2$ level and the H- α peak and plateau fluorescence decay rates (Section 2.2.1 (v), (vi) and (vii)). Nicholson (1983) performed a sensitive Thomson scattering experiment to investigate the existence of a non-Maxwellian electron velocity distribution, and found no evidence of non-Maxwellian behaviour in the electron distribution for energies $< 5 \text{ eV}$. However, use was then made of the ELENDF code (Morgan, 1979) which calculates a self consistent electron velocity distribution given as input data the electron and ion temperatures and densities and atomic level populations. The output from this code indicated that a bump may be formed on the tail of the electron distribution, at energies $> 7 \text{ eV}$, caused by the overpopulation of $n=2$ with respect to LTE. However transitions between excited states are largely unaffected by this bump as it is far from the threshold of transitions of interest and contains very few electrons. Using the collision rates calculated from the ELENDF

code, output from the CR code produced excited state populations which differed by less than 10% from those found by integrating the cross sections over a Maxwellian.

2.2.3 Summary of the Results of the Alternative Rate Formalism. In general, the output from the CR code using the ionization and collision rates of Vriens and Smeets gives a better description of the fluorescence scattering experimental data to date than the output using the rates of Johnson.

Use of the collision rates calculated from the semi-empirical formulae of Vriens and Smeets in the CR code gives good agreement with the experimental data for:

- i) the $n=4$ collision-induced peak enhancement for both the pure hydrogen and 5% hydrogen in helium plasma,
- ii) the $n=4$ collision-induced fluorescence risetime (in particular for $n_e \leq 2 \cdot 10^{14} \text{ cm}^{-3}$),
- iii) the $n=3$ fluorescence transient decay for $n_e \leq 5 \cdot 10^{14} \text{ cm}^{-3}$,
- iv) the $n=2$ population decay whilst pumping the Balmer- α transition, for $n_e \geq 7.5 \cdot 10^{14} \text{ cm}^{-3}$.

There is however no significant difference in the CR code output for:

- i) the pre-laser (quiescent) $n=2$ population density,
- ii) the Balmer- α fluorescence plateau enhancement.

The output from the CR code when employing the rates of Vriens and Smeets are in worse agreement with the experimental data than the output using Johnson's rates for:

1) the Balmer- α fluorescence peak enhancement,

ii) the $n=2$ population decay whilst pumping the Balmer- α transition for $n_e \leq 5 \cdot 10^{14} \text{ cm}^{-3}$.

In the case of the collisionally induced $n=4$ plateau enhancement, the CR code output for the rates of Vriens and Smeets is nearer to the observed values than the output for the rates of Johnson, but is not as good a description of the data than when the rates of Johnson are reduced by a factor of 10.

However, the anomalously high $n=2$ density (for $n_e \geq 3 \cdot 10^{14} \text{ cm}^{-3}$) and the $n=3$ and 4 plateau enhancements observed cannot be explained by a reduction in the accepted collision rates alone. In these cases, an extra rate feeding the $n=2$ level could be responsible and possible mechanisms for such a rate are discussed in the following sections. Note that both the $n=3$ and 4 peak enhancements are described fairly well by the CR code suggesting that these levels could also be overpopulated with respect to their steady state values, to the same degree as the $n=2$ is overpopulated.

2.3 The Possibility of an Extra Rate into the $n=2$ and/or $n=3$ Level.

For an extra rate into the $n=2$ and/or 3 level to be significant, it must be of the same order of magnitude or larger than the main populating collision rate into that level. Tables 2.4 and 2.5 show the main populating processes in the plasma and the associated rates for an electron temperature of 1 eV (equivalent to a delay time of 20 μs in the recombination phase). The relevant electron collision rates of both Johnson and Vriens and Smeets (using the corresponding steady state densities from the CR code) are also tabulated for comparison. After the electron collision rates, the largest reaction

**TABLE 2.4: The Main Processes Populating the n=2 Level
in the 70 cm Z-Pinch.**

Process	Rate into level $T_e = 1 \text{ eV}$ ($\text{cm}^{-3} \text{s}^{-1}$)	Reference
Collisional ^a excitation $\text{H}(n=1) + e \rightarrow \text{H}(n=2) + e$	$5.3\text{E-}13 n_e n_1 = 3.9\text{E}19$	1
Collisional ^a de-excitation $\text{H}(n=3) + e \rightarrow \text{H}(n=2) + e$	$8.41\text{E-}8 n_e n_3 = 1.9\text{E}19$	1
Radiative decay ^a $\text{H}(n=3) \rightarrow \text{H}(n=2) + h\nu$	$4.4\text{E}7 n_1 = 4.0\text{E}18$	2
Dissociative recombination $\text{H}_2 + e \rightarrow \text{H}(n=2) + \text{H}$	$3.61\text{E-}9 n_e n_{\text{H}_2} = 2.3\text{E}17$	3
Electronic recombination $\text{H}^+ + e \rightarrow \text{H}(n=2) + h\nu$	$2.8\text{E-}14 n_+ n_e = 1.7\text{E}17$	4
Mutual neutralization $\text{H}^+ + \text{H}^- \rightarrow \text{H}(n=2) + \text{H}$	$3.0\text{E-}9 n_{\text{H}^-} n_e = 1.5\text{E}17$	5

^aThe ground state $n=1$ density = $3\text{E}16 \text{ cm}^{-3}$ is from the filling pressure, the excited level populations are from the CR code steady state solution (for the rates of Vriens and Smeets).

1 Vriens and Smeets (1980).

2 Wiese et al (1966).

3 Vogler and Dunn (1975).

4 Takayanagi and Suzuki (1976).

5 Bates and Lewis (1955).

**TABLE 2.5: The Main Processes Populating the n=3 Level
in the 70 cm Z-Pinch.**

Process	Rate into level $T_e = 1 \text{ eV}$ ($\text{cm}^{-3} \text{s}^{-1}$)	Reference
Collisional ^a excitation $\text{H}(n=2) + e \rightarrow \text{H}(n=3) + e$	$2.9\text{E-}8 n_e n_2 = 4.9\text{E}19$	1
Collisional ^a de-excitation $\text{H}(n=4) + e \rightarrow \text{H}(n=3) + e$	$5.5\text{E-}7 n_e n_4 = 9.9\text{E}19$	1
Radiative decay ^a $\text{H}(n=4) \rightarrow \text{H}(n=3) + h\nu$	$9.0\text{E}6 n_1 = 6.6\text{E}17$	2
Dissociative recombination $\text{H}_2 + e \rightarrow \text{H}(n=3) + \text{H}$	$4.6\text{E-}8 n_e n_{\text{H}} = 2.9\text{E}18$	3 4 5
Electronic recombination $\text{H}^+ + e \rightarrow \text{H}(n=3) + h\nu$	$1.8\text{E-}14 n_+ n_e = 7.2\text{E}16$	6
Mutual neutralization $\text{H}^+ + \text{H}^- \rightarrow \text{H}(n=3) + \text{H}$	$3.0\text{E-}8 n_{\text{H}^+} n_{\text{H}^-} = 1.5\text{E}18$	7

^aThe ground state $n=1$ density = $3\text{E}16 \text{ cm}^{-3}$ is calculated from the filling pressure and the excited level populations are from the CR code steady state solution (for the rates of Vriens and Smeets).

- 1 Vriens and Smeets (1980) 2 Wiese et al (1966)
 3 Peart and Dolder (1973) 4 Phaneuf and Dolder (1973)
 5 Volger and Dunn (1975) 6 Takayanagi and Suzuki (1978)
 7 Bates and Lewis (1955).

rates populating both levels are dissociative recombination of H_2^+ and mutual neutralization of H^- . Note that the rate into the $n=3$ level is approximately an order of magnitude larger than that into the $n=2$ level. The rates for the H_2^+ and H^- ions were calculated assuming LTE (Saha Boltzmann) densities for these ions. Charge neutrality was also been assumed such that $n_+ = n_e$. Although the electron collision rates at 1 eV appear to be larger than those rates due to other species, specific collision rates may actually be smaller than those given by Vriens and Smeets (as discussed in the previous section) in which case the rates into the $n=2$ and 3 levels from other species may be important.

Reactions such as charge transfer and inelastic scattering between protons and neutral hydrogen atoms, atom-atom excitation and ionization between two neutral atoms only compete with electron collision processes at KeV energies and have thus been neglected. Reactions involving H_2 molecules which populate the $n=2$ or 3 levels directly are also negligible because of the low temperature of the plasma. Dissociation of H_2 molecules by electron impact for example has a threshold of 4 eV (Takayanagi and Suzuki, 1978).

2.3.1 H^- Ions. There is considerable evidence that anomalously high densities of H^- ions exist in low temperature plasmas. Nicolopoulos et al (1977) observed 100 times more H^- in a low pressure, low density (10^9 - 10^{10} cm^{-3}) hydrogen plasma than computed from the relevant mechanisms involved in H^- production and loss. They also found the H^- fraction was high, being up to 35% of the electron density in their plasma. Bacal and Hamilton (1979) observed a nonlinear dependence of the H^- density upon the electron density but pointed out that for $n_e > 10^{10}$ cm^{-3} , the H^- population would be dominated by neutralization and that the H^- density would then be expected to increase linearly with n_e . Other references such as Bacal et al (1981) and Hiskes et al (1982) also

discuss the origin of the high densities of H^- ions found in these low density plasmas.

1) H^- Ion Density. In order to understand the effect that the presence of H^- would have on the $n=2$ and 3 population densities, it is necessary to estimate the density of the H^- ions in the plasma. From Drawin and Felenbok (1965) the LTE density of the H^- ion n_{H^-} is given by

$$n_{H^-} = n_0 n_e \left(\frac{2\pi m_e K}{h^2} \right)^{-3/2} \frac{U_{H^0}}{U_{H^-}} \frac{1}{2} \frac{1}{T^{3/2}} \exp(E/KT) \quad (2-1)$$

where U_{H^0} and U_{H^-} are the partition functions of atomic hydrogen and H^- respectively and E is the ionization potential of neutral hydrogen. All other symbols have their usual meanings. (N.B. The units of this equation are C.G.S.) The assumption of LTE is an approximation to the present plasma as it is known that the $n=2$ atomic level population is overpopulated with respect LTE with the ground state (see Figure 2.2), and therefore steady state calculations were also performed.

Under the plasma conditions investigated, the main rates producing H^- are:



The rate coefficient for the first process ranges from 1 to $3 \cdot 10^{15} \text{ cm}^3 \text{ s}^{-1}$ across the range of plasma conditions studied (Takayanagi and Suzuki, 1978) and for the second process from 7 to $8 \cdot 10^{-31} \text{ cm}^6 \text{ s}^{-1}$ (same reference). The latter process is only important at the higher electron densities. The main process depleting the H^- ion density is mutual neutralization which populates the $n=2$ and 3 levels:



The rate coefficients for the above process range from 3

to $2.5 \times 10^{-9} \text{ cm}^3 \text{ s}^{-1}$ into the $n=2$ state and from 3 to 3.9×10^{-8} into the $n=3$ state for the range of plasma conditions studied here. The rate of change of the H^- population is thus given by

$$\frac{dn_{\text{H}^-}}{dt} = C_1 n_1 n_e + C_2 n_1 n_e^2 - C_{\text{N}2} n_+ n_{\text{H}^-} - C_{\text{N}3} n_+ n_{\text{H}^-} \quad (2-5)$$

where C_1 and C_2 are the populating rate coefficients of processes (2-2) and (2-3) respectively and $C_{\text{N}2}$ and $C_{\text{N}3}$ are the depopulating rate coefficients into the $n=2$ and 3 levels from process (2-4), and n_1 is the ground state density. Assuming charge neutrality ($n_+ = n_e$), the steady state solution for the H^- population is

$$n_{\text{H}^-} = n_1 \frac{(C_1 + C_2 n_e)}{C_{\text{N}2} + C_{\text{N}3}} \quad (2-6)$$

Table 2.6 shows the H^- densities calculated for the plasma conditions of interest for both the LTE and steady-state calculations. The steady state populations are approximately an order of magnitude smaller than the LTE populations.

Delay time (us)	n_e ($\times 10^{14}$)	n_{H^-} (LTE) ($\times 10^{10}$)	n_{H^-} (SS) ($\times 10^9$)
20	24.6	1.9	2.9
30	10.7	3.8	2.5
40	4.6	3.2	1.9
50	2.0	5.0	1.5
60	0.9	8.5	0.7

TABLE 2.6: The Calculated LTE and Steady-State H^- Density.

Note that although the LTE population increases with decreasing temperature and density, the steady state population decreases.

2.3.2 H_2^+ Molecules. Dissociative recombination of H_2^+ populates the excited states thus:



Peart and Dolder (1973) and Vogler and Dunn (1975) measured the cross section of the process into the $n=2$ excited state (using D_2^+ ions) and concluded that the rate into the $n=2$ level was only a small contribution to the total rate. Phaneuf et al (1975) measured the cross section for the formation of the $n=4$ level (in Deuterium) and came to the same conclusion. Vogler and Dunn suggested that the $D(n=3)$ state was the most likely excited state to be populated by this process. This was assumed to be the case for H_2^+ dissociation also, in order to calculate the rates into the $n=2$ and 3 levels from this process.

2.3.2 1) H_2^+ Ion Density. To calculate the density of the H_2^+ ions in LTE, it is necessary to first calculate the LTE density of H_2 molecules in the plasma. Using the molecular Saha equation ($H_2 \rightleftharpoons H + H$), the LTE density of H_2 is given by:

$$n_{H_2^+} = (n_{H_1})^2 \left(\frac{2\pi K T}{h^2} \frac{m_p}{2} \right)^{-3/2} \frac{U_{H_2}}{4} \exp(D/KT) \quad (2-8)$$

Where n_1 is the ground state hydrogen density, U_{H_2} is the partition function of H_2 and D is the dissociation energy of the molecule (4.48 eV). (N.B. the units of this equation are C.G.S.) The expression for the molecular partition function in Mayer and Mayer (1940) and tabulated molecular constants from Huber and Herzberg (1979) were used to calculate the U_{H_2} . The ordinary Saha equation can then be used to calculate the H_2^+ density:

$$\frac{n_{H_2^+}}{n_e} = \frac{n_{H_2}}{n_e} \left(\frac{2\pi m_e K T}{h^2} \right)^{3/2} 2 \exp(-I/KT) \quad (2-9)$$

where I is the ionisation potential of H_2 (15.43 eV) and

all other symbols have their usual meanings. (N.B. The units of this equation are also C.G.S.) The assumption has been made that U_{H_2} is approximately equal to $U_{H_2^+}$ (Nightingale 1982, Spirit 1984).

In order to calculate the steady state H_2^+ density the following populating and depopulating processes were taken into account, since they were found to be the dominant rates under the plasma conditions studied:



The rate coefficient for the process populating H_2^+ (2-10) ranges from 0.3 to $2.5 \times 10^{-17} \text{ cm}^3\text{s}^{-1}$ and for the two depopulating rates from 3.6 to 4.2×10^{-8} and 0.9 to $1.2 \times 10^{-9} \text{ cm}^3\text{s}^{-1}$ respectively (Takayanagi and Suzuki, 1978) for the plasma conditions investigated.

Denoting the populating rate coefficient (process 2-10) by C_1 and the two depopulating rate coefficients (processes 2-11 and 2-12) by C_2 and C_3 and assuming that all the H_2^+ ions produced by process (2-10) are in the ground state (an overestimate), the steady state condition gives the following expression for the H_2^+ density;

$$n_{H_2^+} = \frac{C_1 n_e n_1}{(C_2 n_e + C_3 n_{H_2})} \quad (2-13)$$

Table 2.7 shows the calculated LTE and steady state population for the H_2^+ ions and the LTE population of the H_2 molecules in the 70 cm z-pinch. The H_2^+ density is a small fraction of the electron density and decreases with decreasing electron temperature and density, in the same way as the discrepancy between the experimental and theoretical $n=2$ density.

TABLE 2.7: The LTE and Steady State Density of H_2^+ and the LTE Density of H_2 Molecules in the 70 cm Pinch.

Delay Time (us)	n_e (10^{14} cm^{-3})	H_2^+ density (LTE)	H_2^+ density (steady state)	H_2 density (LTE)
20	24.6	2.56E10	1.73E8	5.25E10
30	10.7	7.10E08	3.50E7	2.14E11
40	4.6	2.56E06	9.68E6	2.08E12
50	2.0	9.82E02	3.03E6	5.43E13
60	0.9	0.01E00	2.96E4	6.52E15

2.3.3 The Effect of H^- and H_2^+ Ions Upon the Atomic Level Populations. The rates into the $n=2, 3$ and 4 levels from the H^- and H_2^+ species were incorporated into the CR code to see if these extra populating processes could account for the anomalies observed from the previous fluorescence scattering measurements. The rise and decay times of the fluorescence signals would not be affected since these are essentially dominated by the inter-level electron excitation/de-excitation rates. However, the absolute level populations, the plateau and possibly the peak enhancement would be affected.

Tables 2.8-10 show the calculated rates into the $n=2,3$ and 4 levels due to the processes involving the H^- and H_2^+ ions compared to the main electron collision rates into these levels. The contribution of the H_2^+ dissociation to the overall level populations is small compared to the contribution of the H^- ions except for the 20 us delay time when they are comparable. The rate into the levels from mutual neutralization of H^- is in general smaller than the main electron collision rates for the earlier delay times, but becomes comparable to or larger than the main collision rates for the later (less dense) times.

TABLE 2.8 to 2.10: The Populating Rates for the n=2,3 and 4 Levels From H^- and H_2^+ Ions Compared to the Main Electron Collision Rates into Each Level.

The rates from the H^- and H_2^+ ions are denoted by the symbols $K_{H^-}n_{H^-}$ and $K_{H_2^+}n_{H_2^+}$ respectively and are quoted for both the LTE and steady state (SS) densities of these ions as previously calculated. The electron collision rates ($K_{ij}n_i/K_{ji}n_j$) of both Johnson (1972) denoted by the letter J, and Vriens and Smeets (1980) denoted by the letters VS are tabulated. All electron densities are quoted in units of 10^{14} cm^{-3} and collision rates are in units of $\text{cm}^{-3}\text{s}^{-1}$.

TABLE 2.8: Rates into the n=2 Level.

Delay (us)	20	30	40	50	60
n_e	24.6	10.7	4.6	2.0	0.9
$K_{12}n_1$ (J)	1.2E20	1.2E18	2.4E15	6.0E11	5.4E06
(VS)	3.9E20	4.9E17	9.6E14	2.3E11	2.0E06
$K_{32}n_3$ (J)	6.3E19	1.1E19	2.8E18	6.4E17	1.2E17
(VS)	1.9E19	3.5E18	8.1E17	1.6E17	2.3E16
$K_{H^-}n_{H^-}$ (LTE)	1.4E17	1.0E17	3.8E16	2.5E16	1.9E16
(SS)	2.1E18	6.7E15	2.3E15	7.5E14	1.6E14
$K_{H_2^+}n_{H_2^+}$ (LTE)	2.3E17	8.0E15	4.6E12	--	--
(SS)	1.6E15	4.0E14	1.7E13	--	--

TABLE 2.9: Rates into the n=3 Level.

Delay (us)	20	30	40	50	60
n_e	24.6	10.7	4.6	2.0	0.9
$K_{23}n_2$ (J)	1.5E20	1.2E19	2.6E18	3.4E17	1.5E16
(VS)	4.9E19	4.6E19	1.1E18	1.5E17	6.3E15
$K_{43}n_4$ (J)	2.5E20	4.0E19	7.2E18	1.4E15	3.0E17
(VS)	9.9E19	1.6E19	2.8E18	5.4E17	1.2E17
K_{H-n_H} (LTE)	1.5E18	1.4E18	5.7E17	4.0E17	3.2E17
(SS)	2.4E17	9.1E16	3.4E16	1.2E16	2.6E15
$K_{H_2}^+n_{H_2}^+$ (LTE)	2.9E18	5.0E16	4.6E13	--	--
(SS)	2.0E16	2.5E15	1.7E14	--	--

TABLE 2.10: Rates into the n=4 Level.

Delay (us)	20	30	40	50	60
n_e	24.0	10.7	4.6	2.0	0.9
$K_{34}n_3$ (J)	3.1E20	3.9E19	6.3E18	7.9E17	6.4E16
(VS)	1.2E20	1.6E19	2.5E18	2.7E17	1.8E16
$K_{54}n_5$ (J)	1.0E21	1.4E20	2.2E19	3.6E18	6.6E17
(VS)	4.4E20	6.1E19	9.3E18	1.5E18	2.9E11
$K_{H_2}^+n_{H_2}^+$ (LTE)	3.3E17	4.5E15	4.6E12	--	--
(SS)	2.3E15	2.2E14	1.7E13	--	--

These extra processes did not make a significant difference to the absolute level populations for both the Johnson and Vriens and Smeets set of rates. The $n=2$ density was increased by approximately 10% and the $n=3$ and 4 levels by less than this. The plateau enhancement was also found to be unaltered because the extra rates populated the $n=3$ as well as the $n=2$ level.

For mutual neutralization to affect the absolute $n=2$ level population, either the H^- density must be larger than that calculated from LTE or the rate coefficients as calculated from Bates and Lewis (1955) must be an underestimate of the true rate. Figure 2.11 shows the $n=2$ density measured by Burgess et al (1980) plus the output from the CR code using the rates of Vriens and Smeets for an arbitrary factor of 10 increase in the rates into the $n=2$ and 3 levels from mutual neutralization. Note that the exact choice of the electron collision rates is not crucial, the steady state densities being relatively insensitive to them. For the range of plasma conditions of interest, the rate from the H^- ions into the $n=2$ and 3 levels is much greater than the H_2^+ contribution. To increase the net rate into the $n=2$ level by a factor of 10, the H_2^+ density would need to be approximately a factor of 50 times higher than its LTE value. Figure 2.11 shows the experimental data from Burgess et al (1980) for the $n=2$ density plus the density predicted by the CR code using the rates of Vriens and Smeets and the extra processes involving i) the LTE densities of H^- and H_2^+ and ii) the LTE density of H_2^+ and 10 times the LTE density of H^- . The experimental data at the higher densities ($n_e > 4.6 \times 10^{14} \text{ cm}^{-3}$) still does not agree with the theory, the measured $n=2$ density being larger by factors of approximately 5. This is an improvement upon the case where no extra rates were incorporated into the CR code, but still fails to describe the experimental data completely.

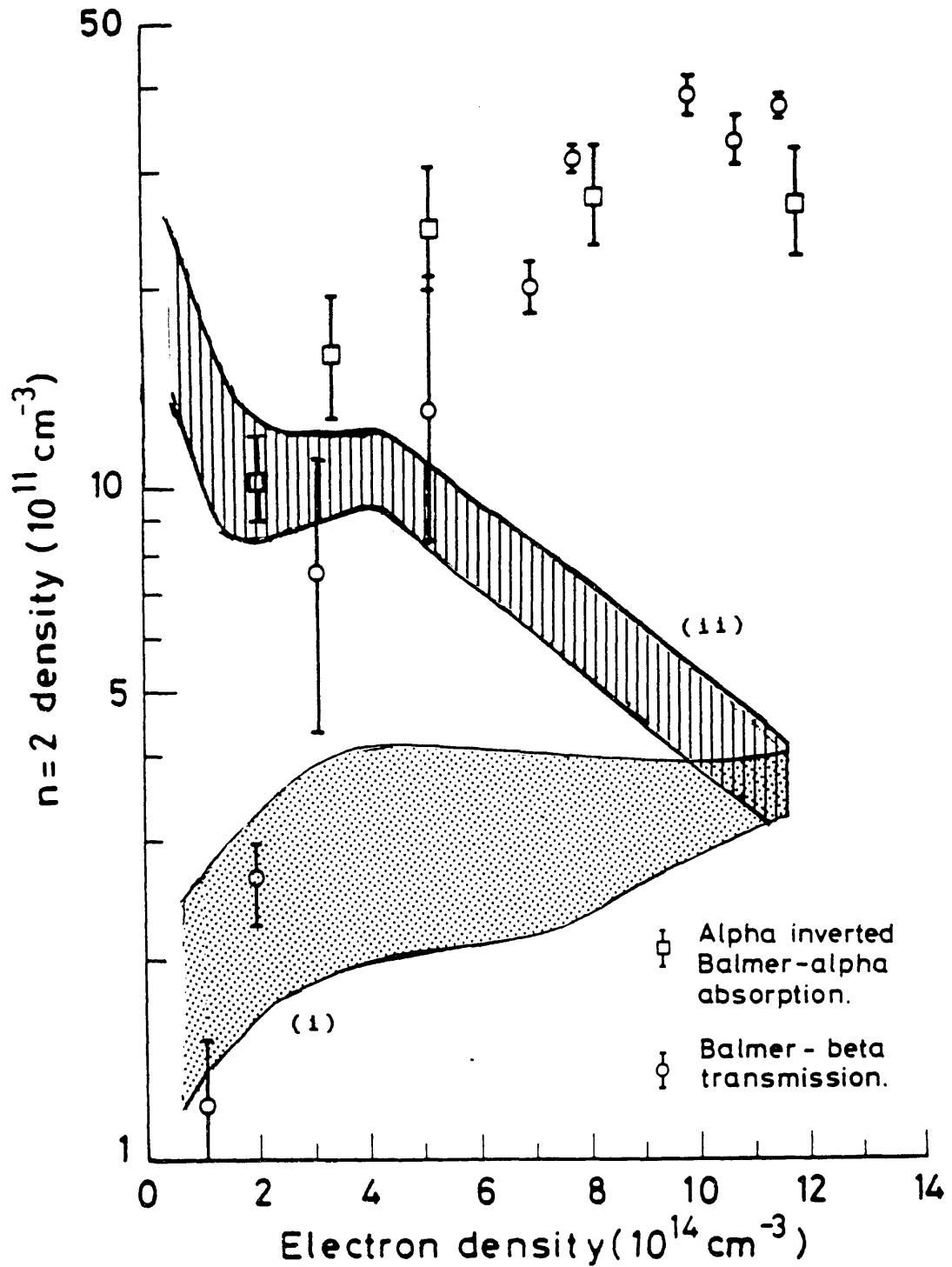


Figure 2.11 The Steady State $n=2$ Level Population as a Function of Electron Density.

The experimental points are reproduced from Burgess et al (1980) and the continuous regions are the theoretical $n=2$ densities as predicted from the CR code for the collision rates of Vriens and Smeets (1980) plus the extra processes into the $n=2,3$ and 4 levels from the presence of:

- i) the LTE densities of both H_2^+ and H^- ions
- ii) the LTE density of H_2^+ and 10 times the LTE density of the H^- .

2.3.4 Summary of the "Extra Rate" Hypothesis.

Processes which populate the hydrogen atomic level populations in the present plasma have been investigated. In general, the rates into the atomic levels from species such as H^- and H_2^+ were comparable to or smaller than the relevant electron collision rates. These extra processes were incorporated into the CR code, assuming LTE densities of these species. The predicted density of $n=2$ was increased by a maximum of less than 10%, and the $n=3$ and 4 densities were increased by less than this. The predicted plateau enhancement of the H- α and sensitised H- β fluorescence signals were unaffected by the inclusion of these processes.

The mutual neutralisation rate was scaled up by an arbitrary factor of ten. This reduced the discrepancy between theory and experiment except for the lower electron densities ($< 2 \times 10^{14} \text{ cm}^{-3}$).

2.4 Conclusions.

The combination of a new set of electron collision rates from Vriens and Smeets (which are lower than Johnson) and the extra processes into the $n=2,3$ and 4 levels from H^- and H_2^+ ions present, plus the addition of error bars on the theory to reflect the uncertainty in the diagnosed temperature, has had the effect of eliminating or reducing some of the discrepancies between theory and previous experimental data. However, some clear discrepancies still exist. They are:

i) the absolute $n=2$ density as measured by Burgess et al (1980), which is higher than theory by an order of magnitude for the higher densities ($n_e > 8 \times 10^{14} \text{ cm}^{-3}$).

ii) the $n=3$ plateau enhancements, which the CR code does not accurately predict across the entire range of plasma conditions investigated. Addition of processes involving

H^- and H_2^+ make no difference.

iii) the collisionally induced $n=4$ plateau enhancements, which are still approximately factors of 1.5 larger than theory even when the rates of Johnson are reduced by a factor of 10. Again, addition of processes involving H^- and H_2^+ makes no difference.

Altering the electron collision rates between the levels en masse has little effect upon the discrepancies above. In the light of these discrepancies the following points need consideration:

i) the H^- density may be higher than the LTE density and/or the mutual neutralisation rate populating the excited levels may actually be larger than calculated from the rate coefficients given by Bates and Lewis (1954). As shown in the previous section, these quantities would need to be more than a factor of 10 larger at the higher electron densities studied to account for the large $n=2$ populations measured.

ii) the H_2^+ density may be higher than the LTE density. The H_2^+ density would have to be at least 50 times larger than LTE to affect the $n=2$ level population. It is unlikely that the dissociative recombination rates are wrong by orders of magnitude having been measured experimentally.

iii) some other process in the plasma may feed the $n=2$ level, such as processes involving H_3^+ and H_2 molecules. No information is available on likely rates feeding the excited levels from H_3^+ and processes involving H_2 molecules are unlikely for several reasons. Firstly, the threshold for dissociation of H_2 which would directly populate the excited levels is approximately 4 eV. Secondly, the $H-\beta$ plateau data from the hydrogen in helium plasma (5% H by volume) shows the same discrepancy with the theory as does the data from the pure

hydrogen plasma. In the H-He plasma, the ground state hydrogen density is a factor of 20 lower than in the pure hydrogen plasma. However, the H_2 density (and thus the rate from H_2 into the $n=2$ level) is a factor of $20^2=400$ times smaller (assuming LTE densities, see Equation 2-8). Thus the discrepancy observed in the H-He plasma is unlikely to have been due to the presence of H_2 molecules.

Chapter Three

BALMER-ALPHA FLUORESCENCE SCATTERING EXPERIMENTS ON A HIGHER TEMPERATURE PLASMA.

In this chapter, a set of Balmer- α fluorescence scattering measurements (obtained in the recombination phase of a similar but slightly hotter z-pinch device than that discussed in Chapter 2) are described. Discrepancies involving the peak and plateau enhancements of the Balmer- α fluorescence signal similar to those obtained employing the previous plasma device are reported. However, the measured decay of the peak enhancement was found to be in agreement with the predictions of the CR code, using rates calculated from the semi-empirical formulae of Vriens and Smeets (1980).

3.1 Aim of the Experiment.

Discrepancies between theory and experiment associated with Balmer-alpha fluorescence scattering have been discussed in Chapter 2. To investigate whether these anomalies were confined to one particular z-pinch it was decided to perform another set of fluorescence scattering measurements employing a similar 70 cm pinch in hydrogen.

3.2 The Z-Pinch.

The z-pinch used in this set of fluorescence measurements was mechanically identical to the z-pinch used for the work by Burgess et al (1980). It consisted of a 70 cm long silica tube of 4.7 cm diameter, charged via a 1 uF capacitor. The current was damped by a 2Ω resistor, having a time constant of 2 us confirmed by the Rogowski coil trace. The experiment was performed at times between 48 and 76 us after peak pinch current, the plasma being essentially current and magnetic field free at these times.

3.2.1 Operating Conditions. The pinch section was first evacuated to less than 10^{-3} Torr and then filled with continuously flowing hydrogen gas at a pressure of 0.45 Torr. The charging voltage was always +20 KV (an input energy of 200 J), and the peak emission was observed to be within 1 us of the peak current flow.

3.2.2 The Electron Density. Electron density measurements by laser interferometry were performed by Kolbe (1985). The electron density n_e at a time t after the peak current flow fitted to an exponential equation of the form:

$$n_e = 33.8 \exp(-t/24.6) \quad (3-1)$$

where n_e has units of 10^{14} cm^{-3} and the time t is in us. The errors associated with the interferometric data were

estimated to be approximately 3%. Figure 3.1 shows the electron density results from interferometry for the pinch as a function of the delay time after the peak current flow. The continuous curve is the best exponential fit to the data (Equation 3.1 above). A comparison with the electron density measurements for the z-pinch described in Chapter 4 (shown in Figure 4.4) shows that for delay times ≥ 30 us (when most of the fluorescence measurements were taken) the pinch described in this chapter is more dense.

3.2.3 Electron Temperature. Thomson scattering in this z-pinch by Nicholson (1983) to determine the electron temperature was carried out under the normal operating conditions. The data was fitted to an exponential equation of the form:

$$T_e = (3.16 \pm 0.3) \exp(-t / (36.4 \pm 2.7)) \quad (3-2)$$

where T_e is the electron density in units of eV, and t is the time after peak current flow in us. The errors associated with the Thomson scattering data were typically 10%. Figure 3.2 shows the electron temperature as a function of the time after peak current flow taken from the data of Nicholson (1983). The continuous curve is the best exponential fit to this data (Equation 3.2). Note that for all delay times considered, the temperature is approximately a factor of 1.7 times higher than for the other z-pinch described in Chapter 4.

In most physical respects, the two z-pinches are identical, the only difference being in the length of the charging cable. In this slightly hotter pinch, the cable is approximately 70 cm long whilst in the other pinch it is much longer measuring approximately 3m. The higher temperature in the first pinch could be due to the lower inductance in the shorter charging cable such that more energy is supplied to the discharge.

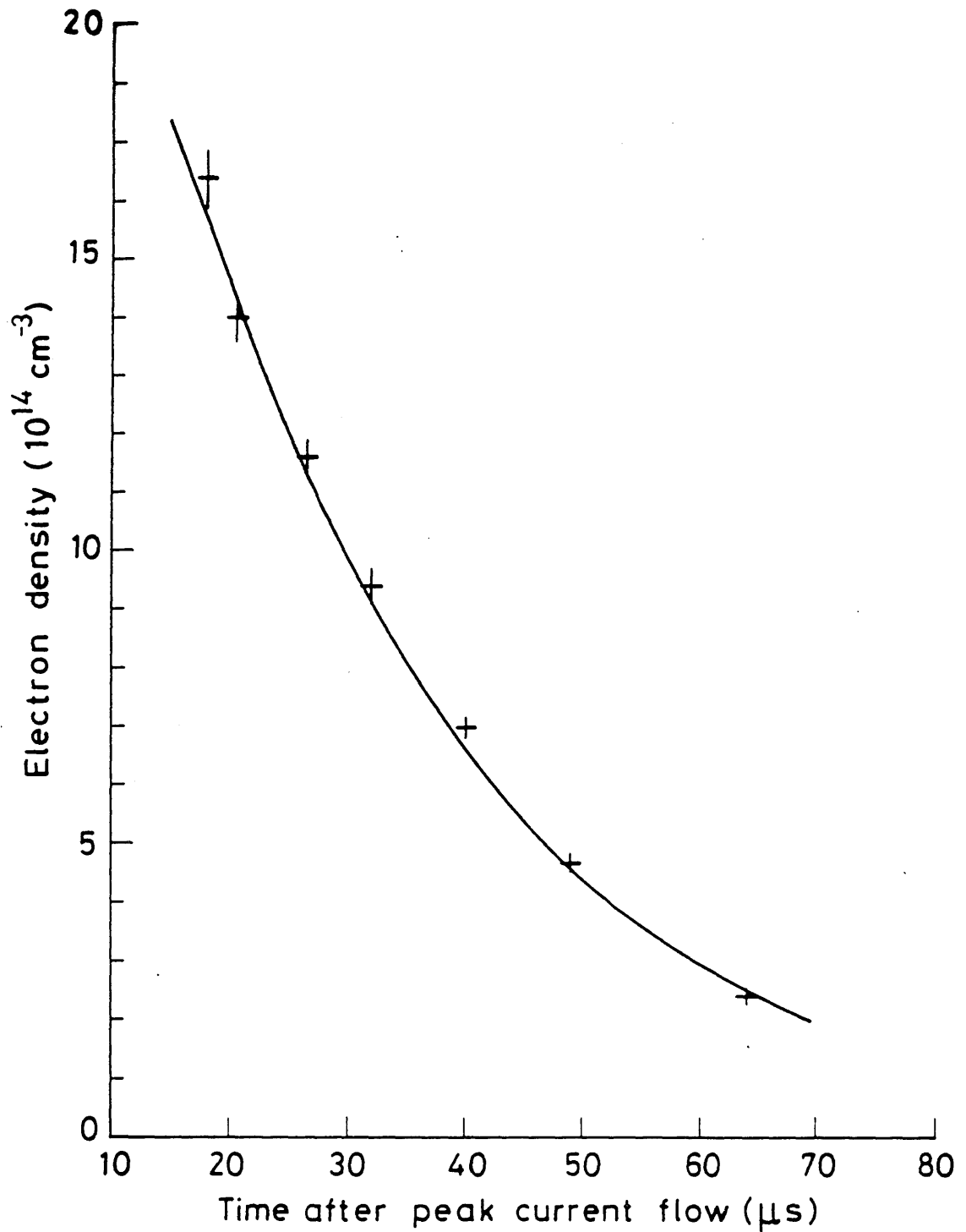


Figure 3.1 Interferometric Electron Density Measurement in 70 cm Z-Pinch.

The data shown is from Kolbe (1985) and the continuous curve represents an exponential fit to the data (Equation 3-1). The filling pressure was 0.45 Torr H_2 and charging voltage 20 KV.

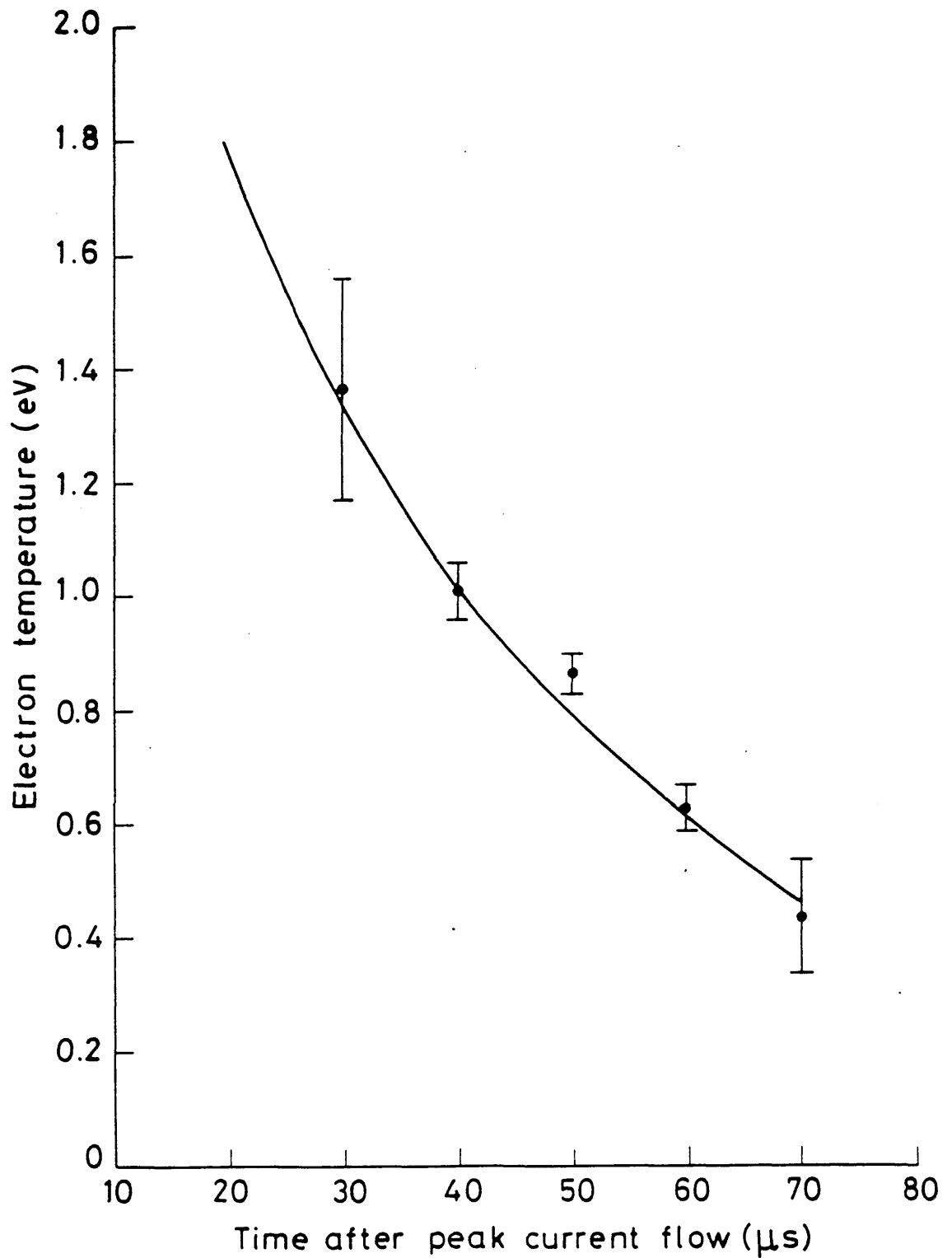


Figure 3.2 Electron Temperature from Thomson Scattering Measurement in 70 cm Z-Pinch.

The data shown is from Nicholson (1983) and the continuous curve represents an exponential fit to the data (Equation 3-1).

3.3 The Coaxial Dye Laser.

3.3.1 Description. The coaxial flashlamp pumped dye laser used was manufactured in-house at Imperial College. It was almost identical to that described by Skinner (1974).

The dye was circulated through a quartz tube lying concentrically within another tube, the 0.5 mm gap being filled with xenon. The discharge across the xenon from a charged 0.3 uF capacitor was initiated by a trigger pulse of 5 KV from a Chelsea Instruments trigger unit, applied to a nitrogen pressurised split-pin type spark gap.

The dye solution for lasing action at the Balmer-alpha (6563 Å) wavelength was a solution of Rhodamine 6G (26 mg l^{-1}) and Cresyl Violet Perchlorate (13 mg l^{-1}) in methanol.

The lasing cavity for the oscillator was formed by two dielectrically coated plane mirrors of reflectivities 100% (the back mirror) and 50% (the output mirror forming the back mirror of the amplifier stage). The amplifier output mirror also had a reflectivity of 50% at the Balmer-alpha wavelength.

The laser was tuned by adjusting the tilt of a 2.5 cm aperture Fabry-Perot of 10 um spacing placed in the oscillator cavity. The divergence of the laser was measured to be 4.5 ± 0.5 mrad.

3.3.2 Laser Operating Conditions. The flashlamps were filled with 60 mTorr of xenon and the charging capacitor always charged to +20 KV. The threshold voltage for successful lasing action was found to be 14.5 KV. The nitrogen pressure in the spark gap was adjusted until the gap was just on the point of breaking down without the trigger pulse being applied (approximately 3 p.s.i). In this way, the jitter in the laser output was reduced to

less than 10 ns.

3.3.3 Laser Characteristics. The Fabry-Perot etalon defined a laser bandwidth of 5 Å, which was checked by scanning across the laser line. Typical laser output energies through the Pockel cell were measured to be 9 mJ in a FWHM pulse duration of 150 ns. This gave a power output of $12 \text{ KWcm}^{-2}\text{Å}^{-1}$, approximately twice as much as that required to saturate the H- α transition at a density of $2.5 \times 10^{15} \text{ cm}^{-3}$ (using the electron collision rates of Johnson (1972) in Equation 1-2). At the later times in the recombination phase less laser power is required to saturate the transition. The 10-90% risetime of the pulse was measured to be $2.0 \pm 0.3 \text{ ns}$ and the e^{-1} fall time approximately 8 ns.

3.4 The Experimental Arrangement.

3.4.1 Description. The dye laser output was passed through an electro-optic switch (a Pockel cell). 30 m of coaxial cable connected between the electrodes defined a laser pulse length of 150 ns. A full description of this particular device can be found in Ward (1983) and Gohil (1982).

The laser output was then passed through the centre of the 70 cm long z-pinch vessel as shown in Figure 3.3 and monitored by a fast photodiode (risetime approximately 1 ns). The photodiode was covered with a 3.0 neutral density filter and a diffuser was used to break up the beam coherence so that speckle did not cause local saturation. The fluorescence signal was viewed side on, perpendicular to the laser beam, as in Figure 3.3, a Rayleigh horn being placed behind the irradiated volume to minimise stray light. The fluorescence signal was dispersed by a 1 metre focal length Hilger Monospek grating spectrometer and detected by an RCA 4836 photomultiplier. The grating used in the monochromator was ruled to 1200 lines mm^{-1} and was blazed for 7500 Å.

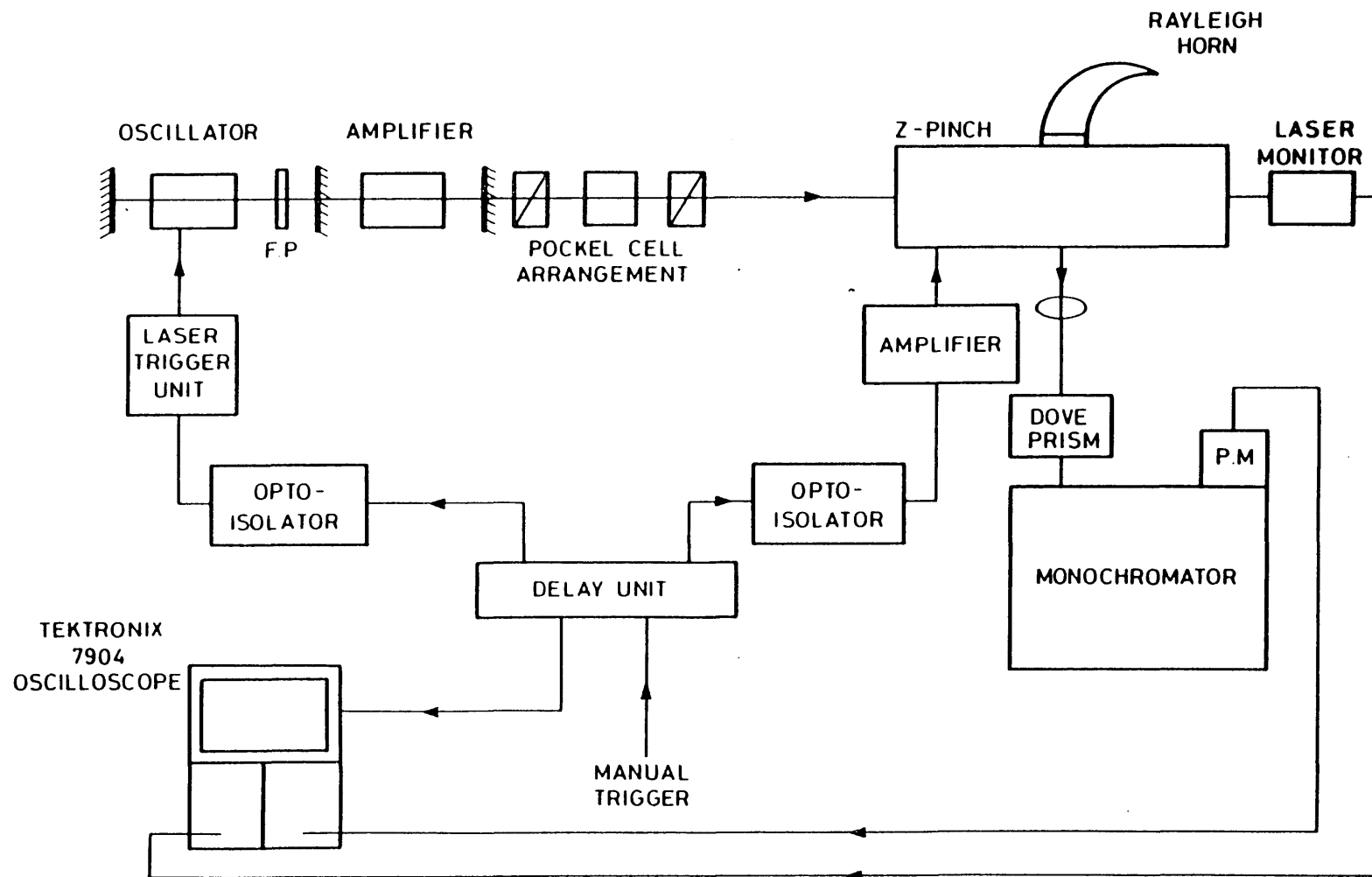


Figure 3.3 Experimental Set-Up for Balmer-alpha Fluorescence Scattering.

The reciprocal dispersion of the detection system was measured to be 9.3 Å/mm using a neon thermaline lamp line at 6328 Å. The photomultiplier was operated at 1500 V. The gain of the tube was typically $\geq 1.5 \times 10^5$ and the quantum efficiency at the H- α wavelength had been measured to be 3.8%.

The output of the photomultiplier was displayed on a Tektronix 7904 oscilloscope with a 7B80 plug-in (1GHz resolution) and a 50 Ω termination impedance and photographed on 612 polaroid film.

3.4.2 Timing. The timing sequence began with a 5V pulse being fed into one channel of a Chelsea Instruments delay unit. The T_0 TTL pulse which was produced synchronously with the triggering pulse was fed into an amplifier which in turn triggered the z-pinch device. The delay unit was set to give the desired delay in the recombination phase and the laser, Pockel cell and oscilloscope were then triggered in sequence so that the resulting fluorescence signal could be viewed. Account was taken of the 2 μ s delay between the plasma device being triggered and the peak of the pinch current, and of the triggering delays of the other pieces of equipment so that the time in the recombination phase at which the fluorescence signal was observed could be accurately known.

3.4.3 Alignment of the Optical System. The rear (100%) mirror of the dye laser was removed and a He-Ne alignment laser was shone through the centre of the quartz tubes, i.e. through the centre of the laser axis. The Pockel cell was then aligned using the following procedure. The two Glan-Taylor polarisers were positioned on axis outside the oscillator/amplifier cavity and rotated relative to each other until a minimum of transmission of the He-Ne was obtained. The Pockel cell was then inserted between the polarisers so that the He-Ne beam passed through the centre of the electro-optic crystal. A

diffuse light source (a tungsten bulb with a tissue in front of it) was placed in front of the first polariser and the Pockel cell carefully rotated until the typical "Maltese Cross" with concentric circles pattern (produced on axis) was centred in the Pockel cell aperture. To align the dye laser into the pinch, one of the polarisers was rotated through 90° so that the He-Ne output was allowed through the Pockel cell, and the position of the plasma vessel was adjusted so that the laser passed through the centre of the vessel.

To align the monochromator, a short focal length lens was used to image the output from a He-Ne laser onto the output slit of the monochromator (i.e. the slit where the photomultiplier tube was to be placed). The He-Ne beam filled the mirrors and grating, and the monochromator was tuned such that the He-Ne light was refocussed onto the input slit. The position of the monochromator was adjusted until the beam intersected the He-Ne beam from the dye laser alignment, in the centre of the pinch. A white card, inclined at 45° to the horizontal and with a black cross drawn on it was inserted into the centre of the plasma vessel at the point of intersection of the two laser beams. The cross was then focussed onto the input slit of the monochromator with an $f/10$ lens. The system was then considered to be optically aligned.

3.4.4 Tuning to Balmer-Alpha. The monochromator was tuned to the Balmer- α wavelength by illuminating the entrance slit with a hydrogen lamp and monitoring the photomultiplier output with a picoammeter. The grating position was adjusted to yield the maximum photomultiplier tube output current. The monochromator was then finely tuned to the peak plasma emission at H- α using a narrow slit width (20 microns).

The dye laser was tuned by carefully directing a little stray laser light onto the entrance slit of the monochromator (which had been previously tuned to the H- α

wavelength) and rotating the intra-cavity Fabry-Perot etalon until the maximum photomultiplier signal was obtained. Care was taken to not saturate the photomultiplier tube. Fine tuning was obtained by rotating the etalon to maximize the actual fluorescence signal from the plasma.

3.5 Results.

Balmer- α fluorescence measurements were obtained for delay times of between 48 and 76 μ s, corresponding to an electron density range of 1.5 to $4.8 \times 10^{14} \text{ cm}^{-3}$ and an electron temperature range of 0.4 to 0.8 eV . All fluorescence traces were displayed on a Tektronix 7904 oscilloscope and recorded on 612 polaroid film, several shots being taken for each delay time.

Plate III 1) shows typical H- α fluorescence traces obtained at three delay times in the recombination phase of the plasma. A typical laser pulse trace is also shown. The fluorescence transient is too fast to be resolved in the case of the 48 μ s data. The laser monitor trace shows one good laser shot and one misfire. Note that the Pockel cell turns on again after the initial termination. This is thought to be due to the reflection of the electrical firing pulse at one electrode.

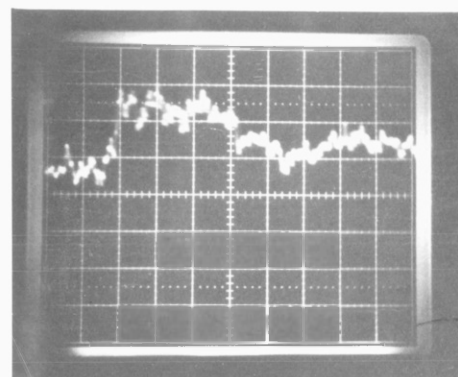
3.5.1 Correction for Optical Depth. The local enhancement observed in the peak and plateau regions of the fluorescence trace must be corrected by the geometric factor g which is the ratio of the irradiated volume to the background volume, which is essentially the ratio of the laser beam diameter to the pinch vessel diameter. The relationship between the local and actual fluorescence enhancement is given by

PLATE III i)

Balmer-Alpha Fluorescence Traces in Hydrogen.

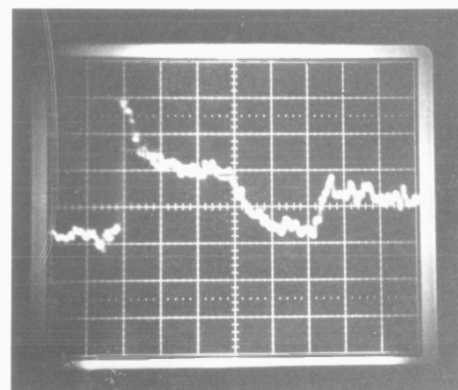
Delay time = 48.00 μ s
 Electron density = $4.80 \cdot 10^{14} \text{ cm}^{-3}$
 Temperature = 0.85 eV
 100 mV/div 50 ns/div

Base line $\longrightarrow$



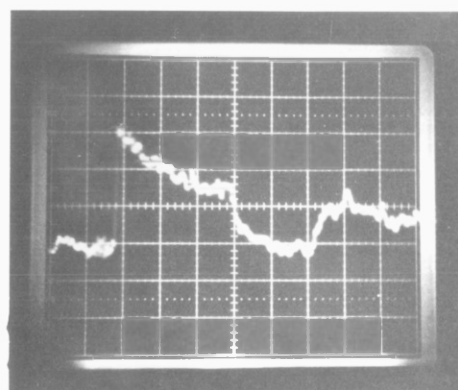
Delay time = 66.00 μ s
 Electron density = $2.31 \cdot 10^{14} \text{ cm}^{-3}$
 Temperature = 0.52 eV

Base line $\longrightarrow$

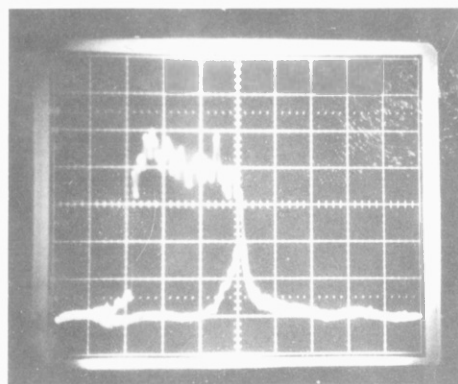


Delay time = 76.00 μ s
 Electron density = $1.54 \cdot 10^{14} \text{ cm}^{-3}$
 Temperature = 0.39 eV

Base line $\longrightarrow$



Laser monitor
 200 mV/div 50 ns/div



$$\text{actual enhancement} = \frac{\text{local enhancement} + 1}{g \times \text{background}}$$

where local enhancement = fluorescence signal above background. In this experiment, $g = (1.03 \pm 0.26) / 4.7 = 0.22 \pm 0.06$. A small correction must also be made for the finite optical depth of the H- α line. Following the analysis presented in Gohil (1982), the correction to the g factor was calculated to be a multiplicative constant of 1.066, giving a final value of $g = 0.24 \pm 0.06$.

3.5.2 Peak Enhancement. Figure 3.4 shows the corrected peak enhancement measured, as a function of electron density. Also shown are the predicted peak values from the CR code for the rates of Johnson (1972). Error bars have been included on the theoretical predictions representing a 10% uncertainty in the diagnosed electron temperature (Nicholson, 1983). (The errors associated with the electron density diagnostic are much smaller, being approximately 3%, also from Nicholson). Each data point is an average value measured over 2 to 4 shots and the error bars are one standard deviation. The horizontal error bars on the data represent a $\pm 4.1\%$ uncertainty in the electron density from an accuracy of ± 1 us in determination of the exact delay time in the recombination phase.

The experimental peak values consistently lie in the lower region of the range of theoretical values. Since the peak enhancement gives (essentially) the ratio of the pre-laser level populations, this suggests that the steady state n=2 level population is overpopulated with respect to the n=3 level population for the range of plasma conditions investigated here. If the lower rates of Vriens and Smeets (1980) are used in the CR code, even higher peak values are obtained which disagree more with the experimental values.

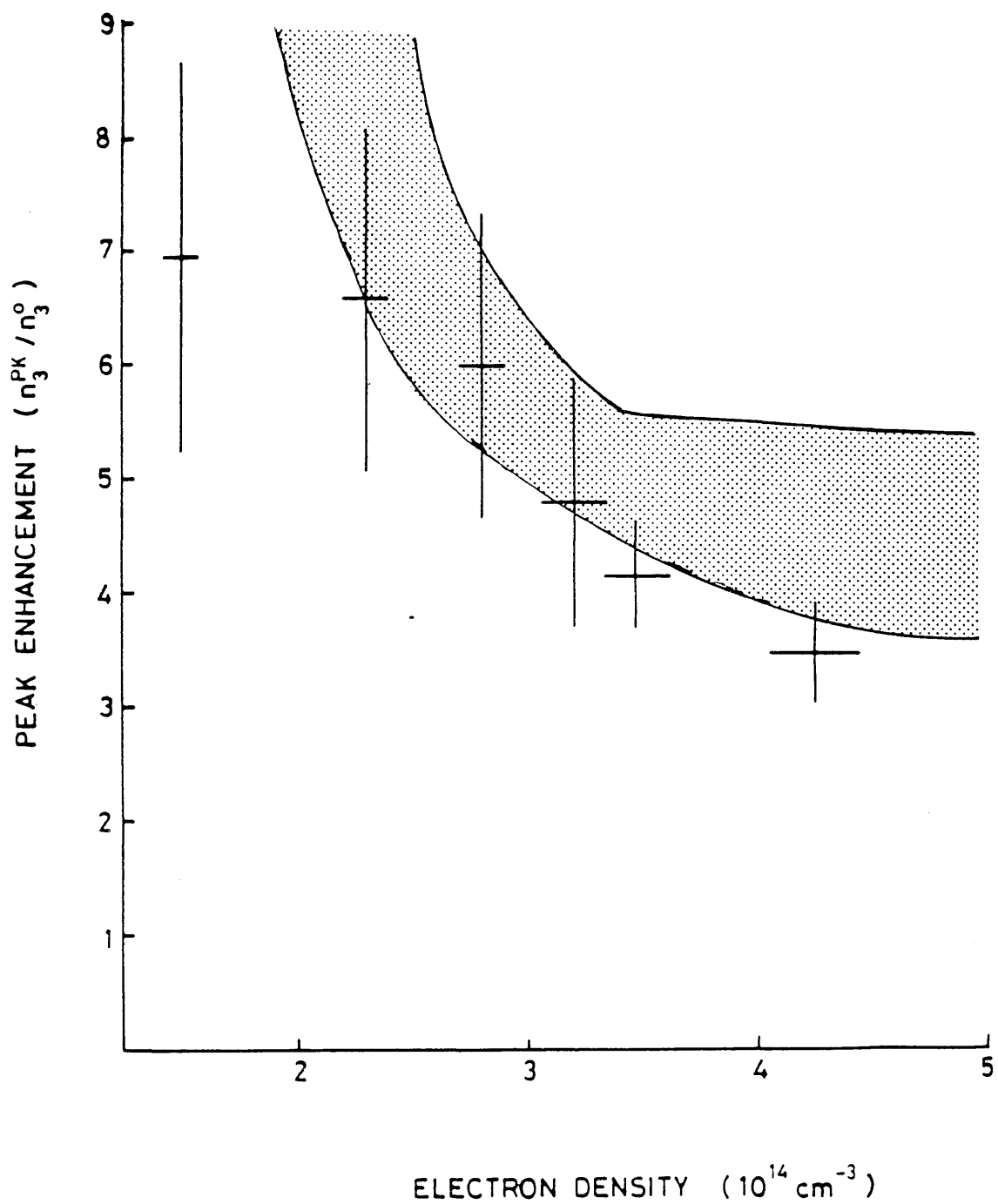


Figure 3.4 The Balmer-alpha Fluorescence Peak Enhancement as a Function of Electron Density.

The continuous regions represent the predicted values from the CR code for the rates of Johnson (1972).

This result is interesting when compared with the results for the peak values measured by Burgess et al (1980) and shown in Figure 2.3, for the same temperature range (from Figure 2.3, this corresponds to $n_e > 5 \cdot 10^{14} \text{ cm}^{-3}$). The peak values also lie below or just within the theoretical values given by Johnson.

3.5.3 H-Alpha Plateau Enhancement. Figure 3.5 shows the measured H- α plateau enhancement as a function of electron density, plus the theory as represented by the output of the CR code using the rates of i) Johnson ii) Vriens and Smeets and iii) Johnson reduced by a factor of 10. The error bars are the same as in Figure 3.3. The experimental data lies above the theory from the rates of Johnson and Vriens and Smeets (except for the single data point at $n_e = 1.5 \cdot 10^{14} \text{ cm}^{-3}$), but is scattered about the predicted values from the CR code for the reduced rates of Johnson. This suggests that the main rates into the $n=2$ and 3 levels which essentially govern the plateau enhancement are actually a factor of 10 smaller than those given by Johnson, or a factor of three smaller than those given by Vriens and Smeets.

The plateau values measured by Burgess et al (1980) and reproduced in Figure 2.5 for the other z-pinch studied over the same range of temperatures investigated here, shows similar discrepancies in that most of the data points lie above the predicted values as given by the CR code for the rates of Johnson and Vriens and Smeets. However, even when the rates of Johnson are reduced by a factor of 10, the CR code output still lies below the measured plateau values unlike the results presented in Figure 3.5.

3.5.4 The H-Alpha Transient (Peak) Decay. Figure 3.6 shows the measured transient decay rate as a function of the electron density, plus the predicted decay rates as given by the CR code using the rates of i) Johnson ii) Vriens and Smeets and iii) Johnson reduced by a factor

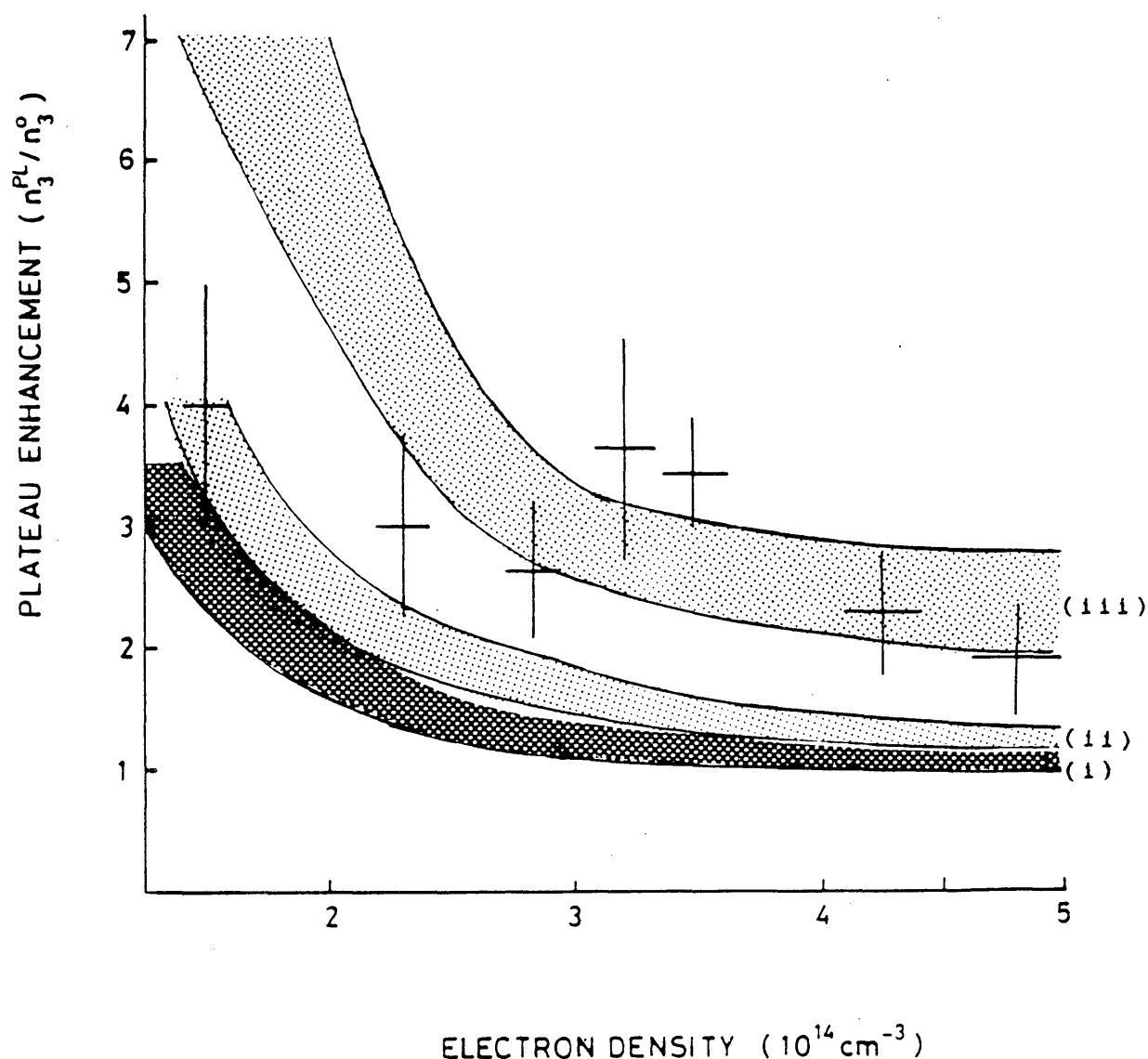


Figure 3.5 The Balmer-alpha Fluorescence Plateau Enhancement as a Function of Electron Density.

The continuous regions represent the predicted values from the CR code for the rates of:

- 1) Johnson (1972)
- 11) Vriens and Smeets (1980)
- 111) Johnson reduced by a factor of 10 (except for the rates to and from the ground state)

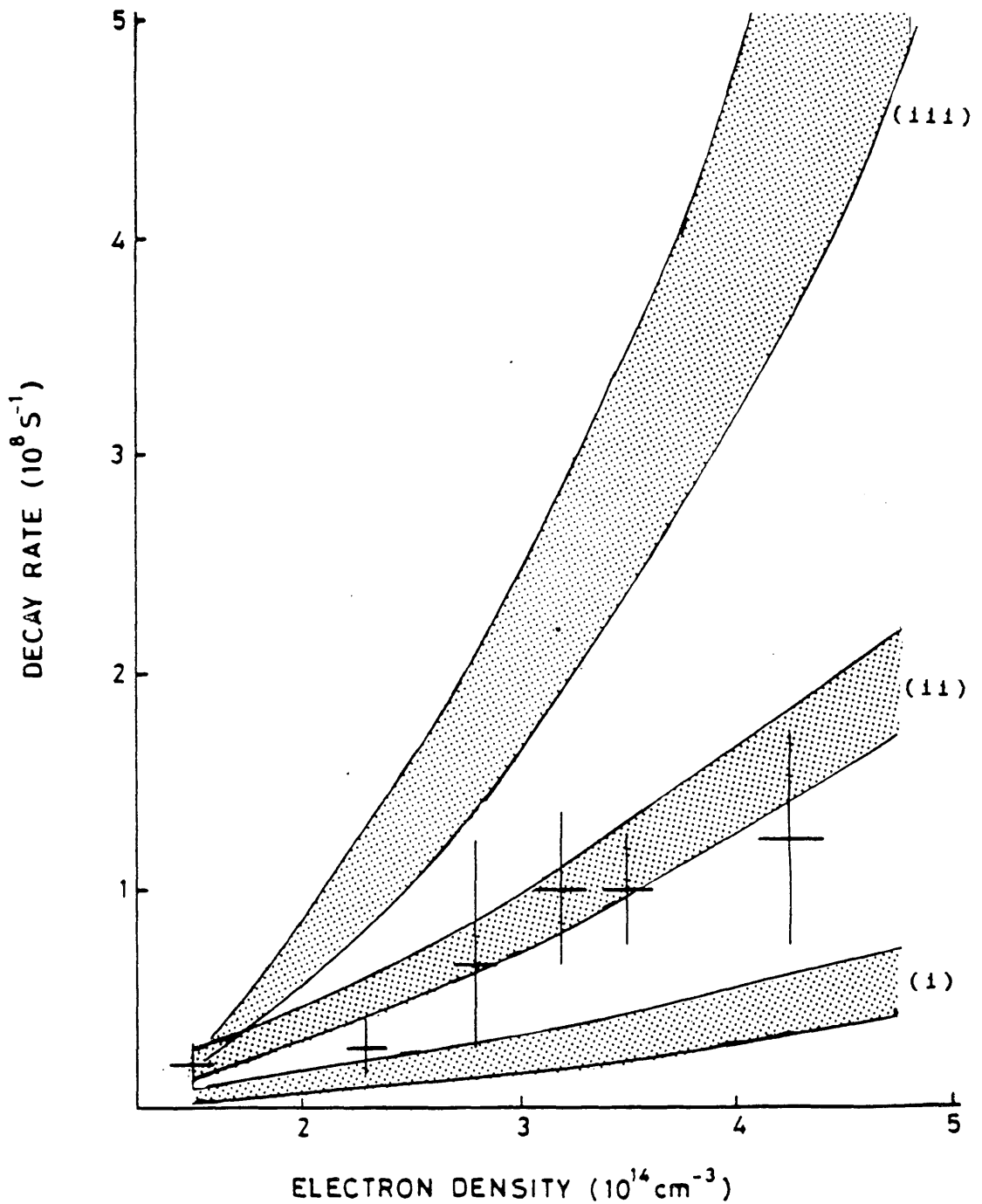


Figure 3.6 The Balmer-alpha Fluorescence Peak e^{-1} Decay Rate as a Function of Electron Density.

The continuous regions represent the theoretical decay rates as predicted from the CR code for the electron collision rates of:

- 1) Johnson (1972)
- 11) Vriens and Smeets (1980).
- 111) Johnson reduced by a factor of 10 (except for those rates associated with the ground state).

of 10. The experimental data agrees well with the output of the CR code for the rates of Vriens and Smeets but lies above the predicted decay rate for the reduced rates of Johnson and below the predicted rate for the unaltered rates of Johnson. The H- α peak decay is governed by the rates out of $n=3$ to the higher levels, essentially the $n=3 \rightarrow 4$ rate K_{34} . The results shown in Figure 3.6 suggest that the K_{34} rate in this plasma for the range of conditions investigated is actually equal to the rate given by Vriens and Smeets, which is typically a factor of 3 smaller than that given by Johnson.

The decay rate for the peak measured by Burgess et al (1980) in the original plasma agrees with the output from the CR code for the reduced rates of Johnson (see Figure 2.8) for $T_e \geq 0.53$ eV, $n_e \geq 4.6 \times 10^{14} \text{ cm}^{-3}$, and with the output from the CR code from the rates of Vriens and Smeets for $T_e < 0.53$ eV, $n_e < 4.6 \times 10^{14} \text{ cm}^{-3}$. This temperature corresponds to an electron density of $2.4 \times 10^{14} \text{ cm}^{-3}$ in the plasma studied in this chapter. At densities greater than this in the present plasma ($T_e > 0.53$ eV) the $n=3 \rightarrow 4$ rate appears to be approximately a factor of 3 larger than that suggested by the decay rate observed in the BMSW plasma.

3.6 Summary and Conclusions.

A series of H- α fluorescence measurements were made using another z-pinch which was slightly hotter and denser than the one used in all previous fluorescence scattering experiments, although it was mechanically identical. Although it was not possible to make an exhaustive set of measurements, similar trends were observed as regards the comparison between the experimentally determined peak and plateau values and the predicted values as given by the CR code output. The peak values agreed quite well with, or lay below, the predicted values from the CR code using the rates of

Johnson (as with the plasma studied by BMSW). Use of the rates of Vriens and Smeets gave predictions in worse agreement with the data.

The measured plateau values were higher compared to the predictions of the CR code for the rates of both Johnson and Vriens and Smeets, as was the case for the BMSW plasma. However, in contrast to the results of BMSW, the plateau enhancement agreed well with the CR code output when the rates of Johnson were reduced by an arbitrary factor of ten.

The high plateau observed by BMSW may be attributed to some process populating the $n=2$ level involving another species. This process may not exist in this plasma.

The $H-\alpha$ peak (transient) decay rate was also measured, and agreed well with the predictions of the CR code, using the rates of Vriens and Smeets. This is also in contrast to the results of BMSW.

Chapter Four

EXPERIMENTAL STUDIES IN THE LOW TEMPERATURE PLASMA.

In this chapter, two experiments are described which were performed to investigate the hypothesis that the presence of H^- ions was responsible for the anomalies associated with the Balmer- α fluorescence scattering in a low temperature hydrogen z-pinch plasma device.

The first experiment was a high power Balmer-alpha fluorescence scattering experiment where the decay of the plateau was investigated. The plateau was observed to decay whilst the laser pulse was still saturating the transition for the earlier delay times investigated, suggesting that H^- does indeed contribute to the plateau enhancement.

Results from the second experiment where the H^- ions were removed by a ruby laser via photodetachment suggested that H^- did not significantly populate the $n=4$ level, or the $n=3$ level at the later delay times. However, to account for the repopulation times observed at the earlier delay times either the H^- density is greater than predicted assuming LTE with the ground state, or the main collision rates out of the $n=3$ level are factors of 2 to 6 lower than given by Vriens and Smeets (1980).

4.1 The Plasma Device.

The set of experiments described in this chapter were carried out (using hydrogen gas) in the recombination phase of the 70 cm long z-pinch, referred to in Chapters 1 and 2 (as distinct from the slightly hotter 70 cm pinch described in Chapter 3). The experiments were performed at times between 20 and 70 μ s after the peak current flow, so that under all conditions of interest the plasma was essentially current and magnetic field free. Under these conditions n_e and T_e lie in the range $0.38 \times 10^{14} \leq n_e \leq 2.5 \times 10^{15} \text{ cm}^{-3}$ and $0.19 \leq T_e \leq 1.03 \text{ eV}$.

A wide range of diagnostic measurements had been previously made on this plasma in the context of spectral line broadening observations (see e.g. Burgess and Cairns (1970, 1971), Burgess and Mahon (1972), Mahon (1973)), far wing Thomson scattering (Nicholson, 1983), and Lyman- α absorption experiments (Baldwin et al, 1984). These measurements, and also unpublished Thomson scattering results by C.C. Smith show the discharge to be homogeneous in density and temperature over a central region of at least 1 cm radial extent, and to within 1 cm of each electrode.

4.1.1 The Operating Conditions. The z-pinch was evacuated to 10^{-3} Torr before being filled to a pressure of 0.45 Torr of continuously flowed hydrogen. The capacitor was always charged to 20 KV (an input energy of 200 J).

4.1.2 Improvements to the Construction of the Z-Pinch.

The z-pinch itself needed to be rebuilt into a more robust and mobile piece of apparatus with a view to moving it from Imperial College to the Rutherford Laboratory for a Lyman- α absorption experiment (Baldwin et al, 1984). The opportunity was taken to add an integral thyristor trigger circuit, mains filters and dump to improve upon the original mechanical construction. Care was taken to keep the original

charging cable at the same length, since this was thought to be connected to the differences in the plasma conditions of the two otherwise identical z-pinch devices discussed in this thesis (see Section 3.2).

A schematic diagram of the spark gap circuitry is shown in Figure 4.1. The 1 μ F capacitor was charged up to the required voltage and the discharge initiated via the trigger circuit. The 2 ohm damping resistor gave a CR time of 2 μ s which was confirmed by observation of the Rogowski trace. The pinch could be triggered via a manual pushbutton, a TTL pulse or an externally generated high voltage trigger. A photograph of the assembled pinch and spark gap is shown in Plate IV 1). The trigger circuit is shown in Figure 4.2. It consisted of a thyristor which, when switched by a TTL pulse produced a pulse of -350 V, with a risetime of approximately 100 ns.

4.1.3 Electron Density Measurement. As a cross check on the constancy of the plasma conditions since the spark gap had been rebuilt, it was decided to remeasure the electron density as a function of the delay time after the peak current by laser interferometry. The interferometric technique of Ashby and Jephcott (1963) was chosen because this method offered greater accuracy than other techniques (such as H- β profile measurements) and could be directly compared with previous interferometry results obtained for this z-pinch.

For a full discussion of the theory of Ashby-Jephcott interferometry, the reader is referred to Ward (1983). The technique involves placing the plasma vessel in an interferometer, and measuring the number of fringes produced when the plasma is formed. The fringes are caused by refractive index variations which are in turn produced by the changes in the electron density of the plasma. Figure 4.3 shows the experimental arrangement of this set of measurements. The visible red and the infra-red line of the He-Ne laser share a common upper

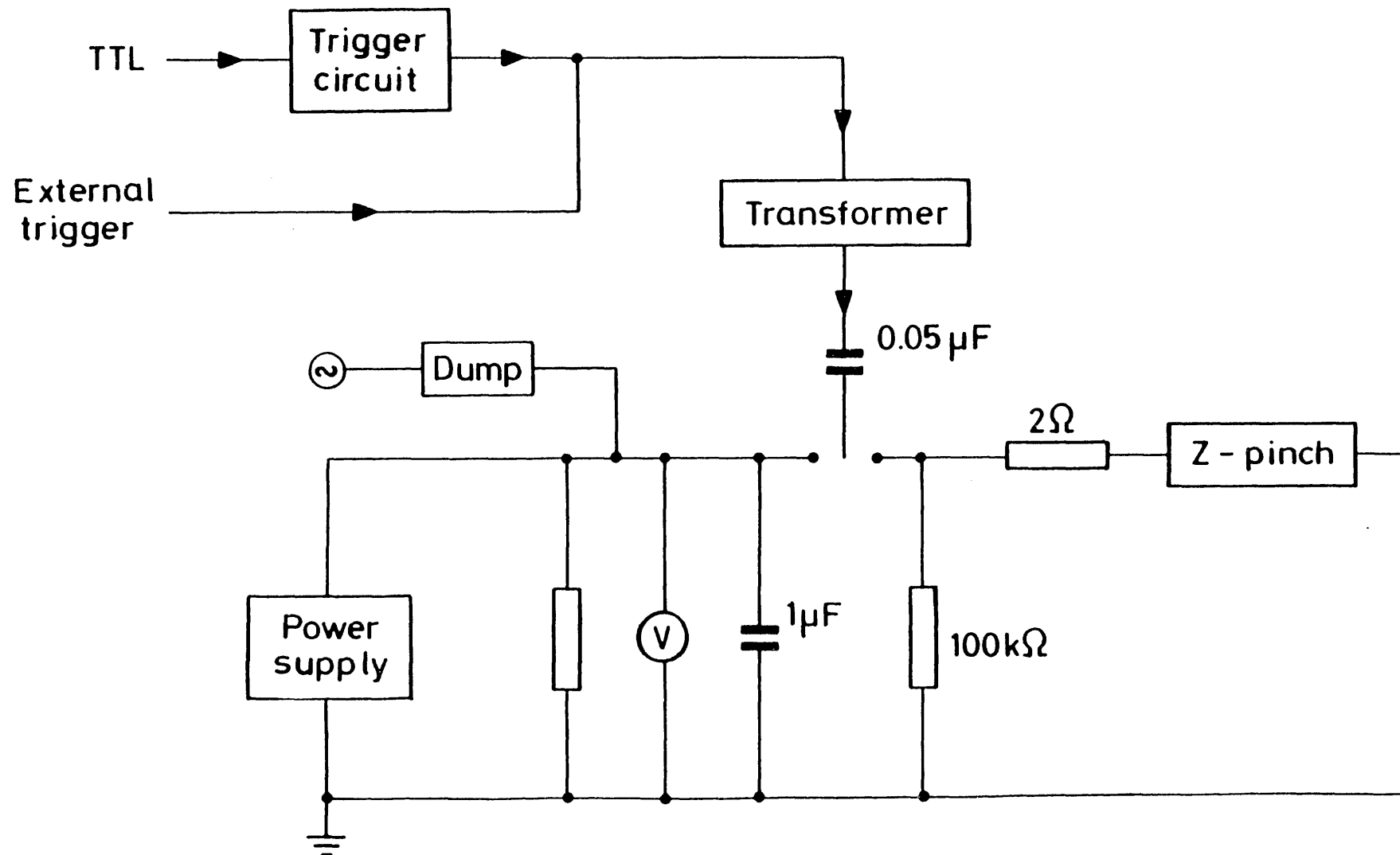


Figure 4.1 A Schematic Diagram of the Z-Pinch Spark Gap Circuit.

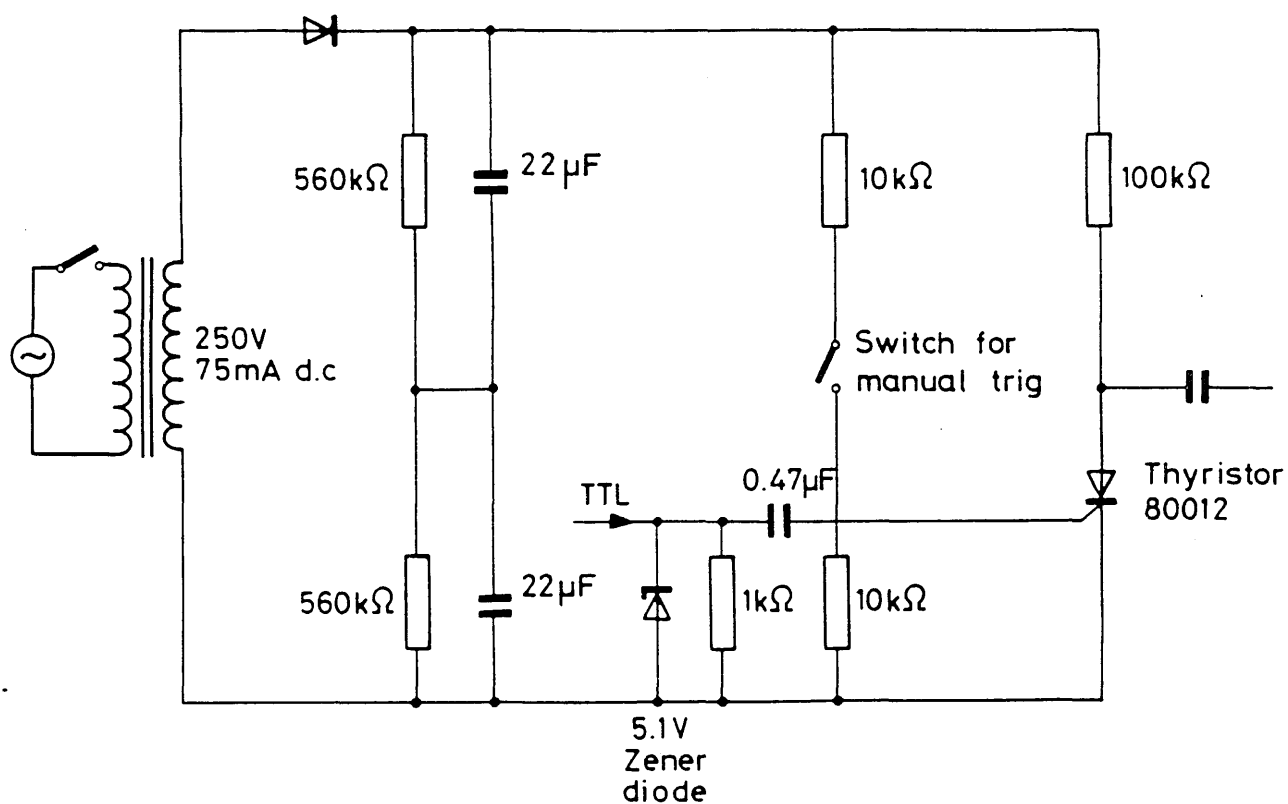


Figure 4.2 The Z-Pinch Trigger Circuit.

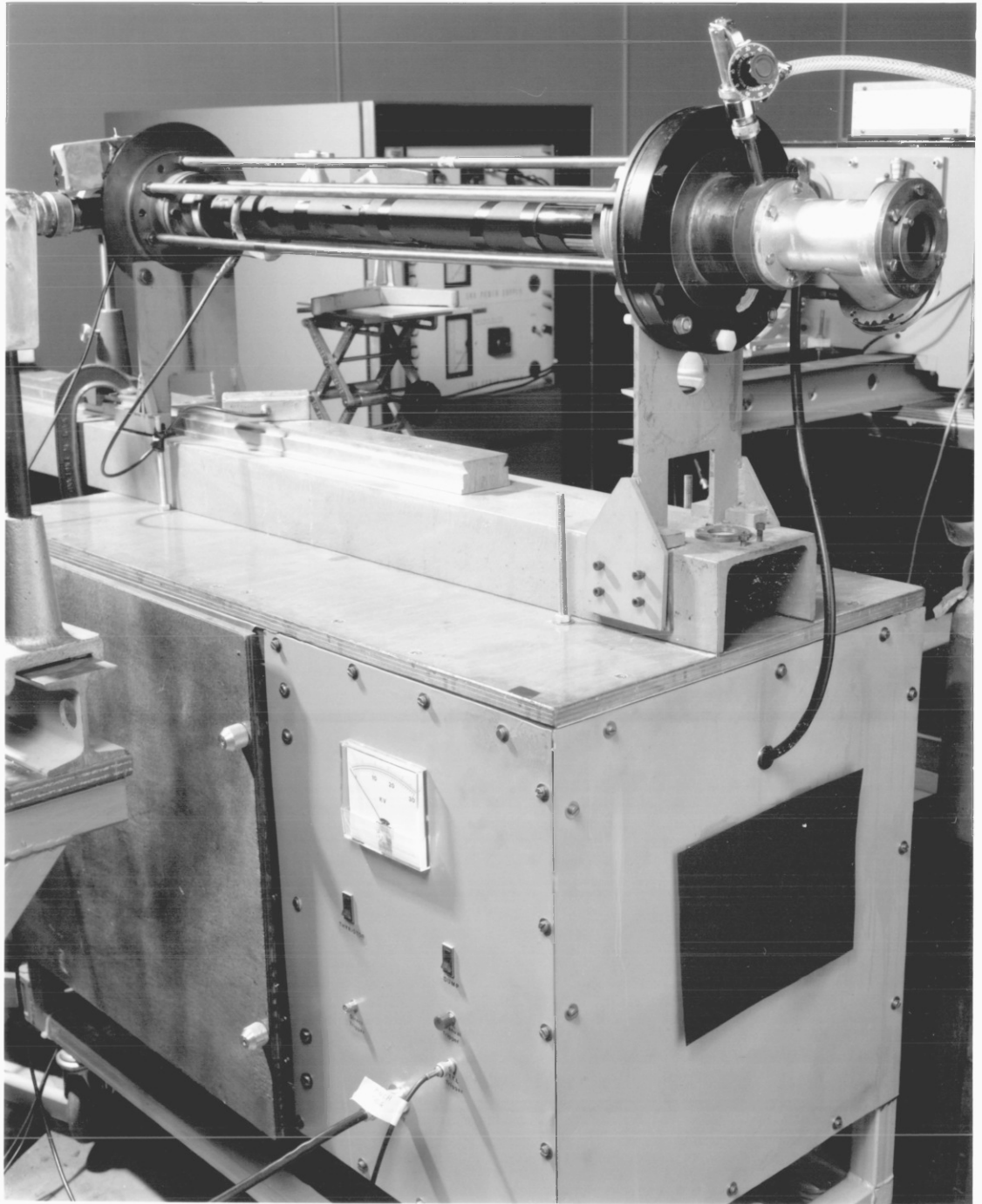


Plate IV 1) The Assembled 70 cm Z-Pinch.

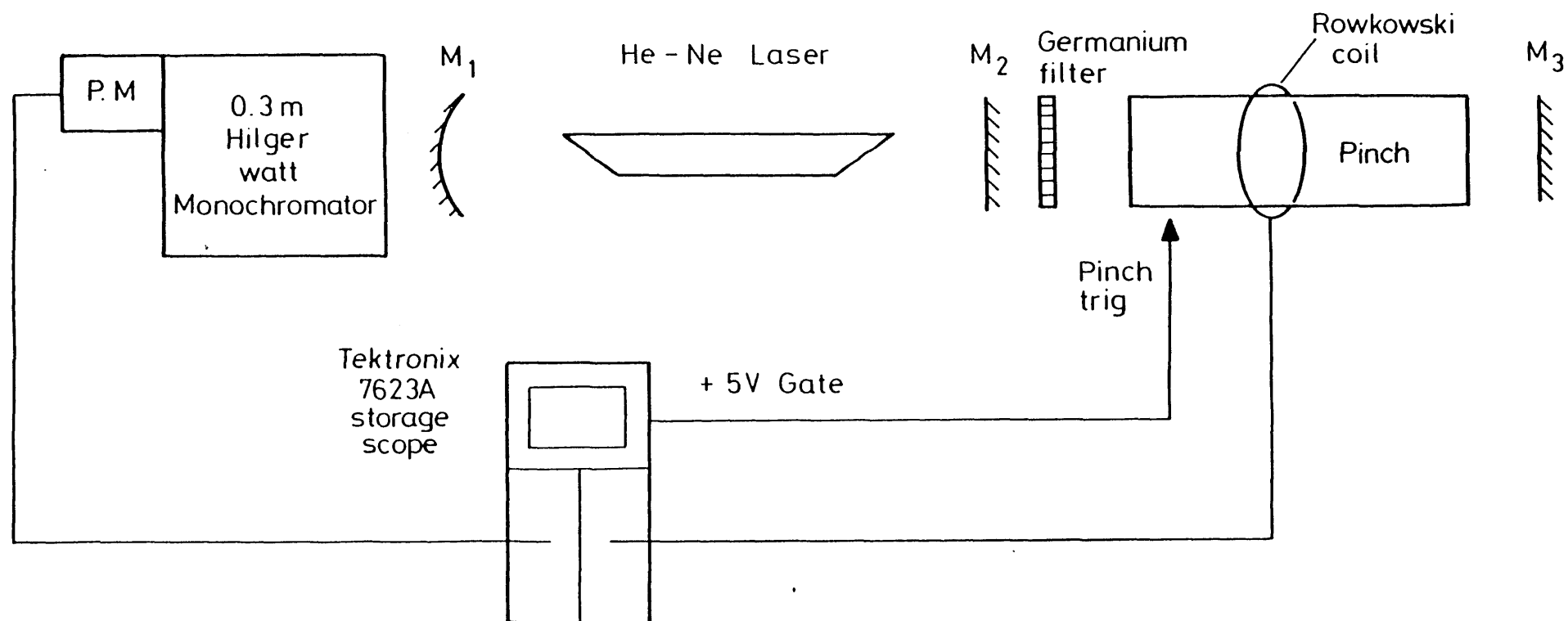


Figure 4.3 Experimental Arrangement for Ashby-Jephcott Interferometer.

state and so modulation of the infra-red line caused by the refractive index changes of the plasma also modulate the visible line. This gives the interferometer the sensitivity of the infra-red line, coupled with the ease of visible light detection. Measuring the position of the fringes observed from the end of the discharge time history (i.e. counting back from a time after peak current flow at which $n_e=0$) gives the electron density as a function of time. Plate IV 11) shows three polaroid photographs of oscilloscope traces obtained in this measurement. The first shows a typical Rogowski coil trace. The second shows a building fringe pattern (due to the apparatus vibration, no plasma) and the third shows the fringes obtained for a typical plasma shot. Note that the building fringes occur over a timescale of milliseconds and therefore do not affect the plasma fringes, which vary over microsecond timescales.

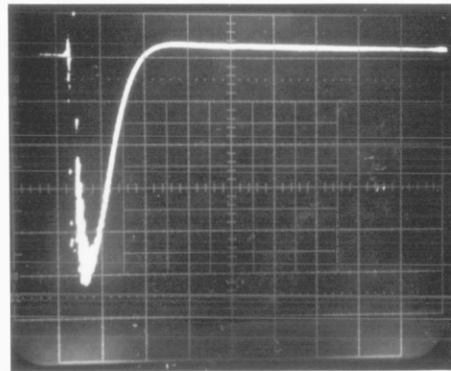
The resulting electron density results are shown as a function of delay time after the peak current in the pinch phase in Figure 4.4. Each data point is the average of at least 8 shots. The continuous curve is an exponential fit to data presented by Burgess et al (1980) and the results presented here agree to within 5% of this and also with the data of Nicholson (1983), Ward (1983) and Mahon (1973). Typically the error bars associated with the electron density measurements were 3%. The rebuilding of the pinch spark gap was therefore deemed not to have affected the plasma conditions.

4.1.4 Electron Temperature. In view of the results presented in Section 4.1.3 above, the results of earlier workers could be adopted for the electron temperature in this plasma. The electron temperature variation with delay time after peak current flow is shown in Figure 4.5. This is taken from Thomson scattering results presented in Burgess et al (1980), and confirmed by further extensive Thomson scattering (Nicholson, 1983). Typical error bars were 10% in this case. The continuous

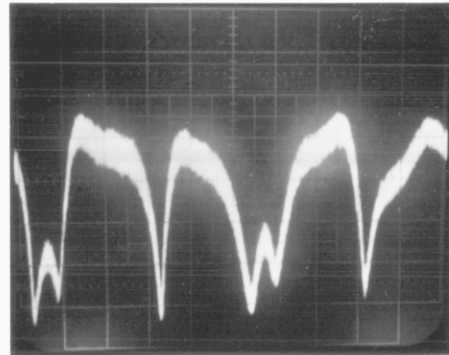
PLATE IV ii)

Laser Interferometry Raw Data

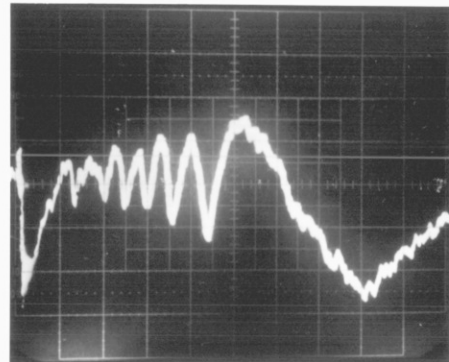
Rogowski coil trace
integrated into 1 M Ω
200 mV/div 5 μ s/div



Typical building fringe trace
(no plasma)
200 mV/div 5 ms/div



Typical plasma shot showing
interference fringes plus
Rogowski trace
200 mV/div 10 μ s/div



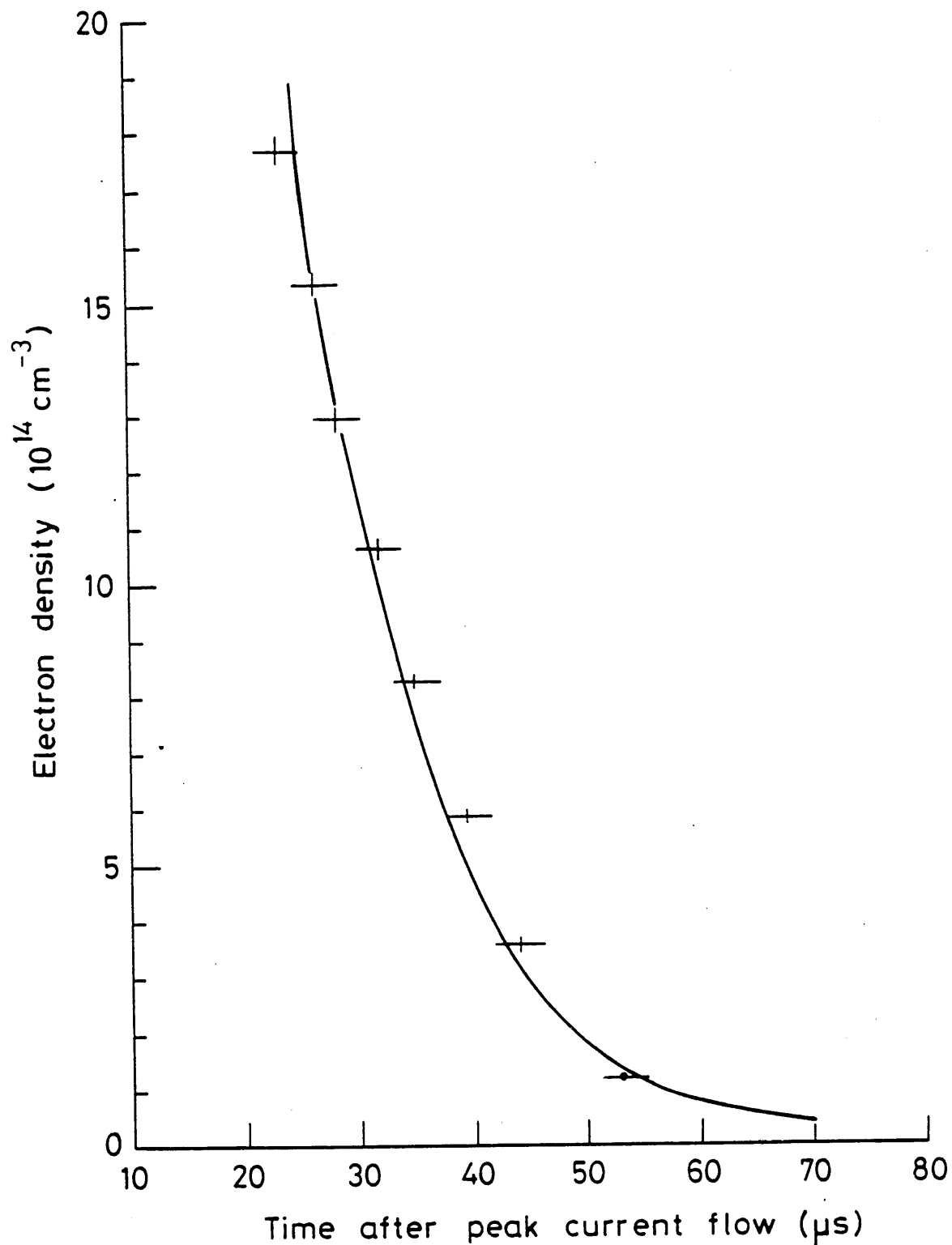


Figure 4.4 Interferometric Electron Density Measurement in 70 cm Z-Pinch.

The crosses on the graph are the data points obtained by the author and the continuous curve represents an exponential fit to data from Burgess et al (1980) for the electron density in the same z-pinch device. The filling pressure was 0.45 Torr H_2 and charging voltage 20 KV in both cases.

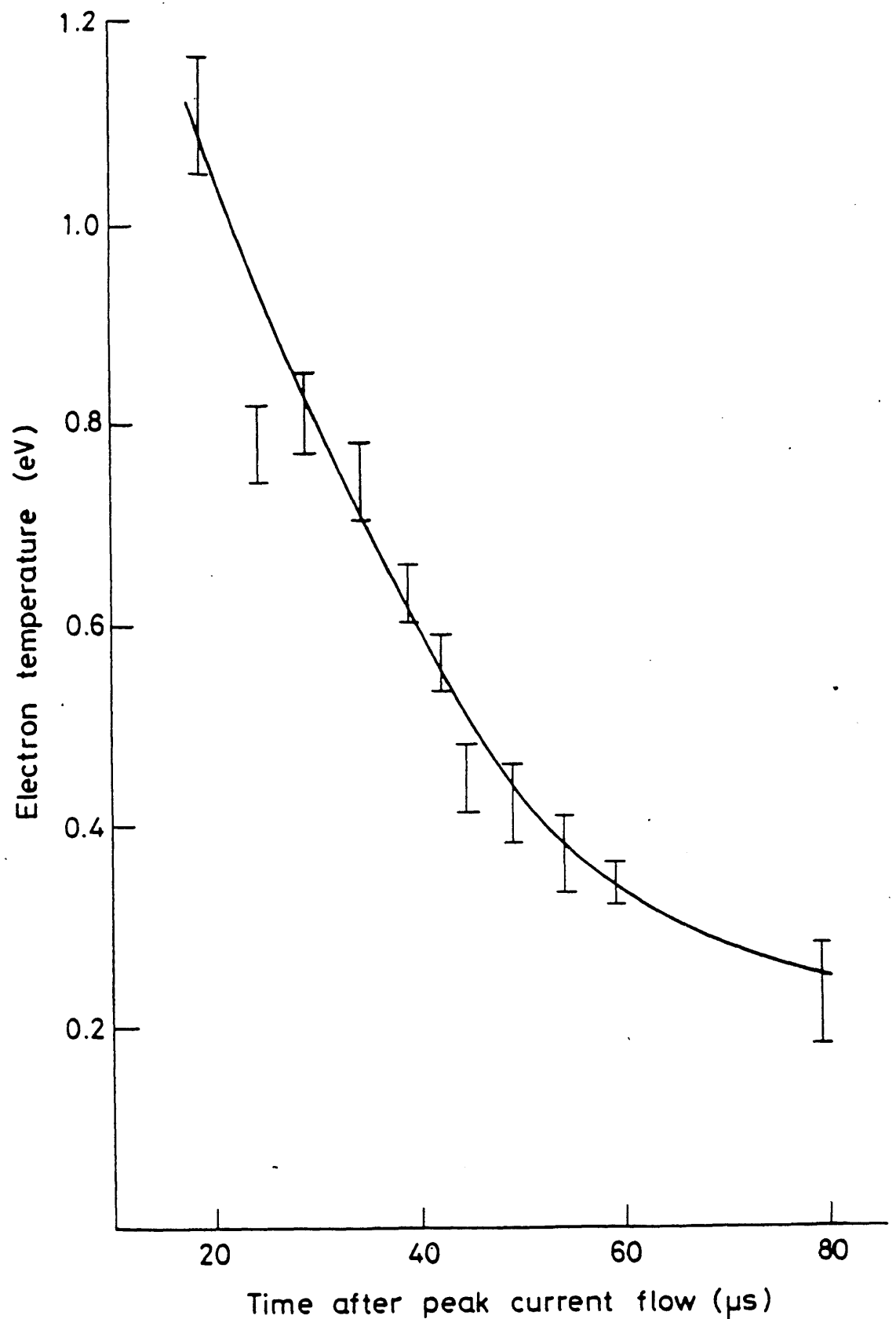


Figure 4.5 Electron Temperature as a Function of Time After Peak Current Flow.

The continuous curve represents an exponential fit to data which has been taken from Thomson scattering results presented by Burgess et al (1980).

curve represents an exponential fit to the data points.

4.2 The High Power Fluorescence Scattering Experiment.

The purpose of this experiment was to remove a significant proportion of the H^- ions presumed present by the use of an intense $H-\alpha$ laser pulse, and to observe the subsequent time history of the characteristic fluorescence plateau to investigate whether it would begin to decay whilst the laser pulse was still saturating the transition. If the plateau was observed to decay prematurely, this could be taken as evidence that the anomalously high plateau is produced by a species that had been removed by the intense laser beam. As has been discussed in Chapter 2, the H^- ion is the species most likely to contribute to the anomalous plateau enhancement.

4.2.1 Overview of the Experimental System.

1) The Dye Laser. The dye laser used in this experiment was the same coaxial flashlamp pumped oscillator/amplifier arrangement used in Chapter 3 on the higher temperature z-pinch, and a more detailed discussion of this particular laser is to be found there. The bandwidth $\Delta\lambda$ of the laser was measured to be $5\text{\AA}$ and the divergence (half angle half power) 4.5 ± 0.5 mrad by focussing the laser through a variable aperture and measuring the power as a function of the aperture diameter. The laser was tuned to the $H-\alpha$ wavelength using a hydrogen lamp reference source and a monochromator as described in Section 3.4.4. The dye solution used was a mixture of Rhodamine 6G (at a concentration of 26 mg/l) and cresyl violet (13 mg/l) dissolved in methanol.

ii) The Laser Output Power. When operated at a charging voltage of 20 KV, the output energy of the laser was typically 95-105 mJ (as measured with an NPL- calibrated calorimeter) at $\lambda=6563$ Å, with a pulse duration of approximately 400 ns (FWHM). With a bandwidth $\Delta\lambda$ of 5 Å this gave approximately 19-21 mJ into the H- α linewidth, corresponding to a power of 50-70 KWcm⁻². Assuming 8% losses due to the input window of the plasma vessel, a power of approximately 46-64 KWcm⁻²Å⁻¹ was expected at the scattering volume, which was more than a factor of 3.5 times greater than the power required theoretically to saturate the H- α transition at the earliest delay time (from Equation (1-2), using Johnson's set of collision rates). For later times, less power would be required for saturation.

Ward (1983) reported powers in the range 17-30 KWcm⁻²Å⁻¹. Although the laser powers obtained in the present work are not orders of magnitude greater than those reported previously, the duration of the laser pulse in the present work is significantly longer than previously used (400 ns FWHM, compared to 150 ns in the previous work which employed a Pockel cell).

iii) Experimental Arrangement. The dye laser output was passed through the centre of the pinch vessel as shown in Figure 4.6 and was monitored by a fast photodiode in the same way as described in Section 3.4.1. The fluorescence signal was dispersed by the same grating spectrometer and detected by an RCA C31034A photomultiplier.

Care was taken to switch the photomultiplier supply voltage slowly up to the operating voltage of the tube (-1400 V) and to ensure that the dark current did not exceed 100 nA, as damage results to the sensitive photocathode when D.C. signals of greater than this are drawn. The dark current was continually monitored with a picoammeter and was typically less than 40 nA. The

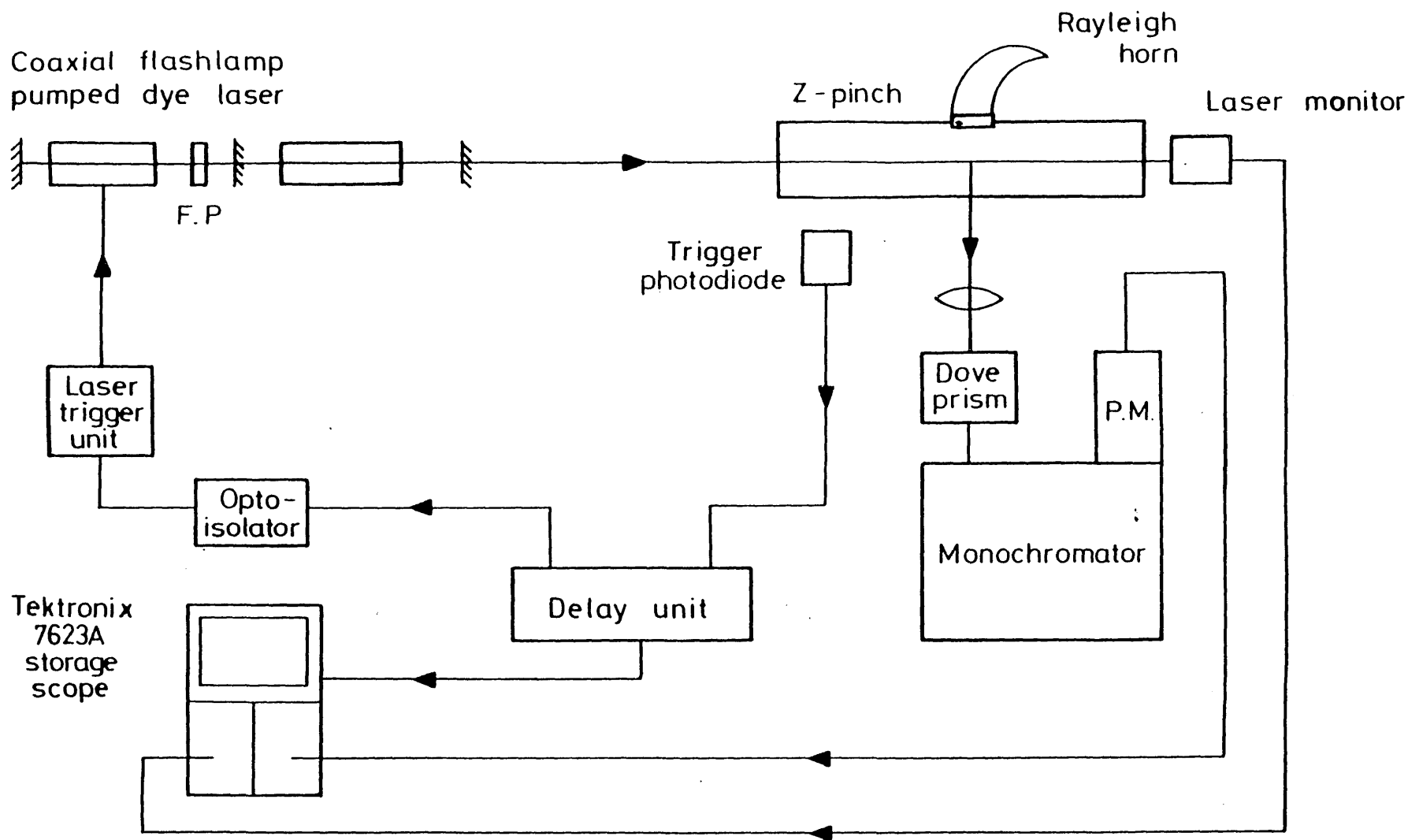


Figure 4.6 Experimental Arrangement for High Power Balmer-Alpha Fluorescence Experiment.

photocathode had a quantum efficiency of almost 15% at $\lambda=6563 \text{ \AA}$ and the gain of the tube was typically 10^6 at the working voltage of -1400 V. The output from the photomultiplier was displayed on a Tektronix 7623A storage scope with a 7A12 plug-in and 50Ω terminating impedance. The whole system was aligned in the way described in Chapter 3, section 3.4.3.

The detection system was checked to ensure linearity by placing neutral density filters in front of the monochromator input slit.

With the high power laser shots used in these experiments, stray laser light swamping the fluorescence signal became a significant problem. The usual unported z-pinch vessel was replaced by a vessel with 1 cm long viewing ports added. These ports were perpendicular to the pinch axis. A Rayleigh horn was placed on one port to form a viewing dump. With the addition of these ports, stray light did not affect the measurements further.

However, Lyman- α emission and absorption studies on the same ported z-pinch section under the same operating conditions showed that cool plasma in these ports produces appreciable end-zone effects (Baldwin, Burgess, Evans, Marangos and Nicholson, to be published). It is not currently possible to quantify the effect of these end zones on the absolute signal levels obtained in the present work. Qualitatively however, the presence of cool end zones in the viewing ports will i) decrease the background level, due to absorption and ii) increase the fluorescence enhancement during the laser pulse as some of the emission from the end zone furthest from the monochromator entrance slit will be viewed as the laser pumping action reduces the opacity of the pumped volume. The actual enhancements observed will therefore be larger than those obtained in previous work using the unported section. In fact, the peak enhancements observed were 20% larger than previous work, and the plateau enhancements

were between 50% and 400% (depending on the plasma conditions) larger than previous work.

iv) Timing. A photodiode was positioned to observe the plasma emission. This diode was connected to an amplifier which produced an output pulse of +5V when triggered, and was sufficiently far away from the plasma to be triggered only by the plasma emission and not by electrical pickup associated with the plasma initiation. The plasma device itself was triggered by an external TTL pulse. The pinch emission triggered the photodiode/amplifier and the TTL pulse produced was fed into a Chelsea Instruments delay unit. This pulse occurred at the peak of the plasma emission, which corresponded to the maximum current flow in the pinch phase (as determined by a Rogowski coil) to within ± 1 μ s. The delay unit was set to give the desired delay in the recombination phase and the output from this unit was used to trigger both the oscilloscope and the laser, the laser being triggered via a fibre-optic link to prevent electrical pickup from affecting the timing chain. The laser itself was placed in a screened box for the same reason. Allowance was made for the delay of 600 ns between the trigger signal arriving at the laser triggering electronics and the initiation of lasing action.

v) Checks for Saturation. For each set of fluorescence scattering measurements made, checks were performed to ensure that saturation of the transition had been achieved. This is particularly important in the present experiment as any observed decay of the fluorescence plateau could be attributed to non-saturation of the transition if these checks were not made.

A 0.3 neutral density filter was placed in the laser beam path. This attenuated the laser by a factor of $10^{0.3}=2$. For each set of measurements made, less than 10% variation in the magnitude and time history of the fluorescence signal was observed, and therefore the laser

was considered to be saturating the transition. On an individual, shot to shot basis, if the fluorescence signal rose to its peak value faster than the laser pulse risetime (approximately 110 ns), then this also indicated that saturation had been attained.

4.2.2 Effect of the Laser on the Plasma. Before discussing the high power fluorescence scattering results, it is necessary to discuss briefly the effect of the laser beam upon the plasma. The power of the H- α laser was sufficient to saturate the n=2 $\rightarrow$ 3 transition and produce the characteristic fluorescence observed. The laser would be expected also to have the following effects.

1) H⁻ Photodetachment. H⁻ has only one bound level with a detachment energy of 0.75 eV, corresponding to a wavelength of 1.66 μ m. If irradiated by photons with energy greater than the photodetachment threshold, the following process occurs:



The cross section for this process $\sigma(\nu)$ was taken to be $3.8 \times 10^{-17} \text{ cm}^2$ at the H- α wavelength (Takayanagi and Suzuki, 1978). The energy of a typical H- α shot was measured to be 100 mJ into a 5 Å bandwidth at the plasma entrance window. Assuming 8% losses at this window, this gives 92 mJ into a FWHM pulse width of 400 ns at the plasma centre. The flux density Φ was therefore calculated to be $5.9 \times 10^{23} \text{ photons s}^{-1} \text{ cm}^{-2}$. The rate R of photodetachment is:

$$R = \Phi \sigma(\nu) \quad (4-2)$$

For the values quoted above, $R = 2.2 \times 10^7 \text{ s}^{-1}$ per H⁻ ion. To calculate the theoretical photodetachment fraction $\Delta n_- / n_-$, Equation 1 from Bacal and Hamilton (1979) was used:

$$\frac{\Delta n_{H-}}{n_{H-}} = 1 - \exp \left\{ - \frac{\phi a(\nu)}{(area) h \nu} \right\} \quad (4-3)$$

where ϕ is the energy of the laser pulse. From the figures already provided and taking an irradiated area of 1.3 cm^2 (the radius of the laser beam being 0.65 cm at the centre of the pinch from the measured divergence) gives $\Delta n_{H-}/n_{H-} = 1$, and so effectively all the H^- present in the plasma can be assumed to be removed during a typical $H-\alpha$ pulse produced by the present laser.

ii) Photoionization from $n=3$. The cross section for photoionization from the $n=3$ level in hydrogen is $3.8 \times 10^{-17} \text{ cm}^2$ from Takayanagi and Suzuki (1978) at 6563 Å. For the range in laser energies quoted and assuming the absorption of the photons results only in photoionization, the total photoionization rate per $n=3$ atom is calculated to be in the range $4.6 \rightarrow 6.1 \times 10^6 \text{ s}^{-1}$.

The stimulated de-excitation rate from $n=3 \rightarrow 2$ is in the range $4 \rightarrow 5 \times 10^{10} \text{ s}^{-1}$ per $n=3$ atom. This rate is approximately 4 orders of magnitude higher than the photoionization rate, and therefore photoionization may be neglected. (Using the previous experimental evidence (discussed in Chapter 2), photoionization may also be neglected compared to the electron collision rates. (However, if the electron collision rate between $n=3 \rightarrow 4$ at the 70 us delay time is equal to that given by Vriens and Smeets ($7.7 \times 10^6 \text{ s}^{-1}$), then photoionization may not be negligible compared to the collision rate.)

The photoionization rate for $n > 3$ falls rapidly since it varies as n^{-5} (Kramers rule) and so photoionization of the higher levels was assumed to be negligible.

iii) Changes in Plasma Conditions. Gohil (1982) investigated the effect of a typical $H-\alpha$ laser pulse on the plasma conditions. He estimated that the maximum change in the electron density due to photoionization would be

2.6% and the maximum change in the temperature to be 5.8%. The effect of the H- α laser on the actual plasma conditions was thus assumed to be negligible. Nicholson (1983) performed Thomson scattering synchronously with a saturating H- α laser pulse and found no change in either the total scattering (proportional to the electron density) or in the $1/e$ width of the scattered spectrum (which gives the electron temperature).

iv) Continuum Absorption. Spirit (1984) showed that the largest contributions to the hydrogen continuum opacity under similar plasma conditions as discussed here were from the hydrogen bound-free, the H $^-$ bound-free and the H $^-$ free-free continuum absorption. The maximum absorption coefficient for the H bound-free continuum was calculated to be $1.5 \times 10^{-5} \text{ m}^{-1}$ for the 20 μs delay time from Thorne (1974), the level populations being taken from the CR code steady state values (the rates of Johnson (1972) were used). The maximum absorption coefficient for the H $^-$ bound-free contribution was calculated to be $7.1 \times 10^{-4} \text{ m}^{-1}$ for the 70 μs delay time from Geltman (1962) assuming Saha-Boltzmann H $^-$ densities. The H $^-$ free-free absorption coefficient was calculated to be $2.0 \times 10^{-4} \text{ m}^{-1}$ at 20 μs from Stilley and Callaway (1970). Absorption of the laser by the continuum for the 35 cm of the hydrogen plasma traversed by the laser before reaching the volume under investigation should therefore have been less than 0.05% at the most.

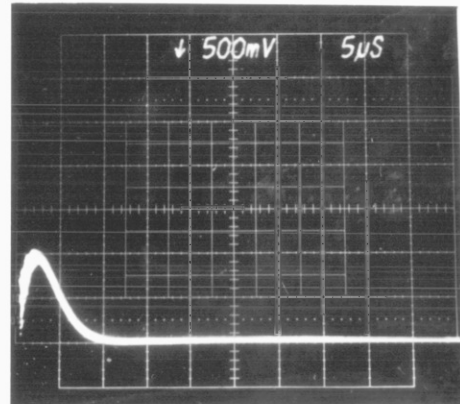
4.2.3 Experimental Results. The results were displayed on a Tektronix 7623A storage scope and photographed using Polaroid 667 film. Data was taken at delay times of 30, 50 and 70 microseconds after the peak pinch current flow.

Plate IV iii) shows the Rogowski coil output (with a 1 M Ω terminating impedance), the H- α plasma emission throughout the duration of the plasma, and a typical laser pulse as detected by the fast photodiode (measured with a 50 Ω terminating impedance). The FWHM of the laser

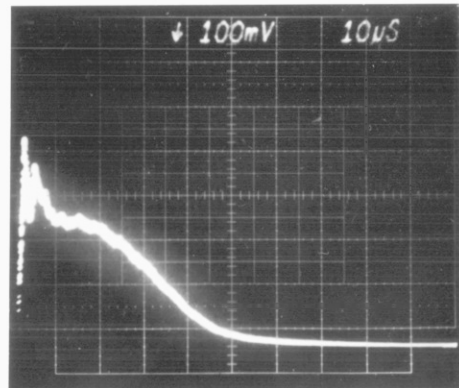
PLATE IV iii)

High Power Laser Fluorescence Data

Rogowski coil trace



Entire pinch emission
at Balmer-Alpha into 50Ω



Laser monitor (photodiode)

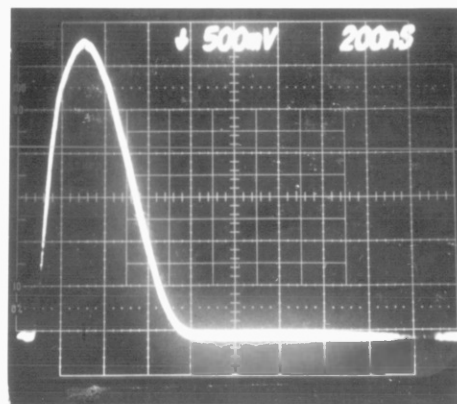
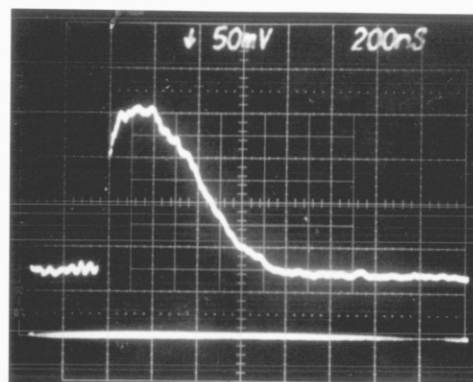


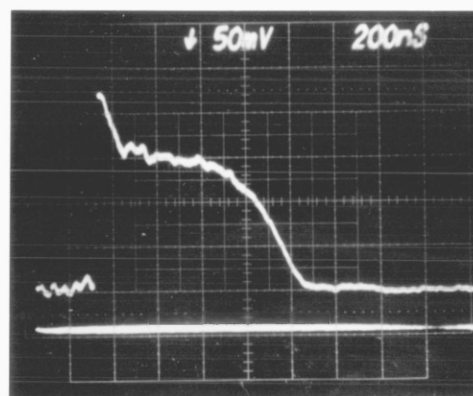
PLATE IV iv)

Balmer-Alpha Fluorescence Traces in Hydrogen

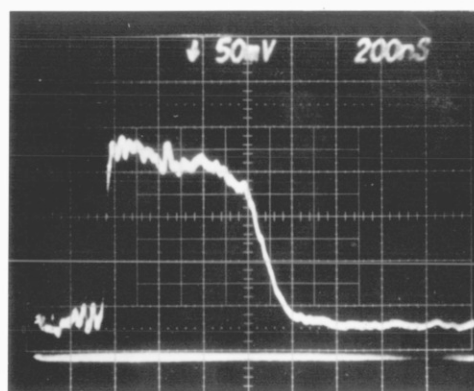
Delay time = 30.00 μ s
 Electron density = $10.70 \cdot 10^{14} \text{ cm}^{-3}$
 Temperature = 0.74 eV



Delay time = 50.00 μ s
 Electron density = $2.02 \cdot 10^{14} \text{ cm}^{-3}$
 Temperature = 0.38 eV



Delay time = 70.00 μ s
 Electron density = $0.38 \cdot 10^{14} \text{ cm}^{-3}$
 Temperature = 0.19 eV



pulse was measured to be 400 ns.

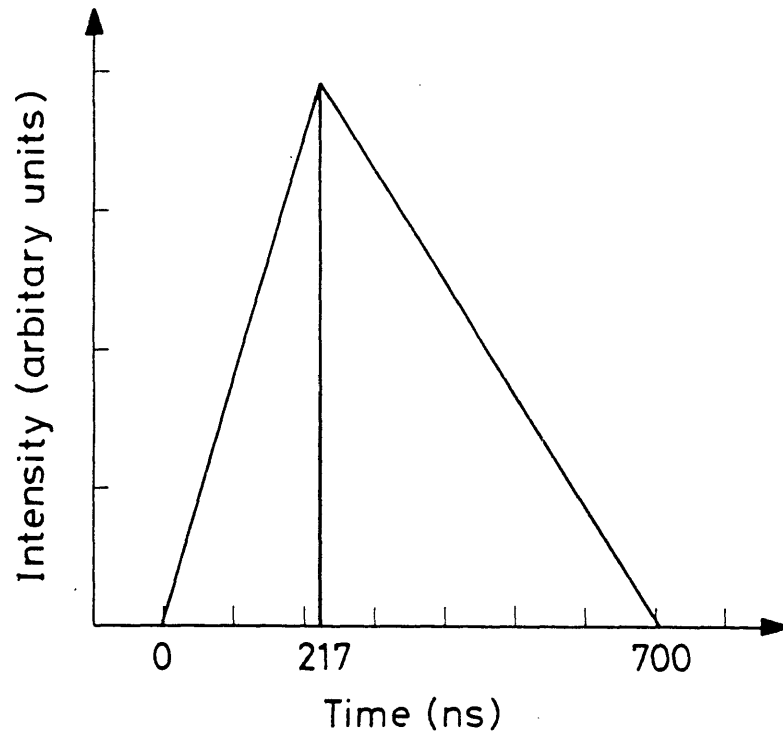
A typical fluorescence trace at each of the three delay times studied is shown in Plate IV iv). Note that the characteristic fluorescence spike is clearly visible in the 50 us delay time trace, but is too fast to be resolved on the timescale used for the 30 us delay time trace. (Burgess et al (1980) measured the transient decay time to be approximately 13 ns at this delay time). It is not clear from the 70 us delay trace whether a transient exists. The CR code predicts a flat topped signal for the corresponding plasma conditions.

4.3 Discussion of the High Power Fluorescence Experimental Results.

4.3.1 Plateau Decay. The laser pulse time history was incorporated in the CR code to study the effect of the varying laser intensity on the fluorescence plateau. The laser pulse shape was approximated to by a triangular shape as shown in Figure 4.7, the peak intensity being approximately 217 ns after laser initiation. Also shown in Figure 4.7 is a tracing of a typical laser pulse recorded using a photodiode with a risetime <1 ns.

The lowest laser energy measured (after allowing for entrance window losses) after an experimental run was 60 mJ. Assuming an irradiated area of 1.3 cm^{-2} (calculated from the laser divergence) this gives a minimum energy of $9.2 \text{ mJcm}^{-2}\text{A}^{-1}$. The expected fluorescence signal was generated from the CR code using this pulse shape and total energy, using the collision rates of Vriens and Smeets.

1) Approximation to Laser Pulse.



ii) Actual Laser Pulse (Photodiode Trace).

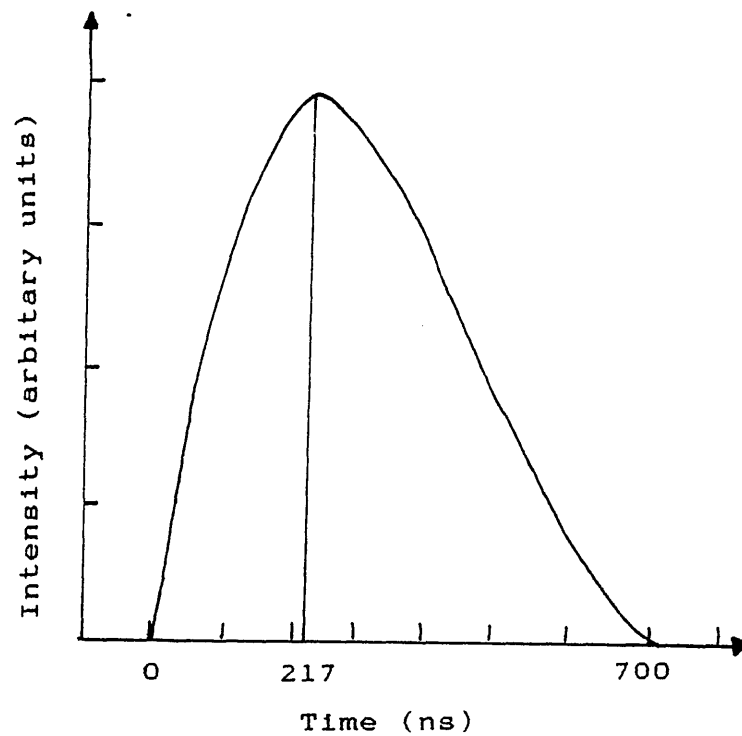


Figure 4.7 The Laser Pulse Time History.

For the analysis, the laser pulse was approximated to be of a triangular shape as shown, the peak in the intensity being approximately 217 ns after laser initialisation. Also shown is an actual laser shot recorded with a fast photodiode (<1 ns risetime) for comparison.

i) 30 us delay. Figure 4.8 shows a typical fluorescence trace obtained at this delay time, and the prediction of the CR code employing a pulse of 60 mJ energy into the laser bandwidth (which was the minimum energy measured after an experimental run) dispersed over the triangular pulse shape discussed above. The CR code output has been normalised to have the same plateau height as the experimental trace, simply so that the time histories of the two traces can be compared easily.

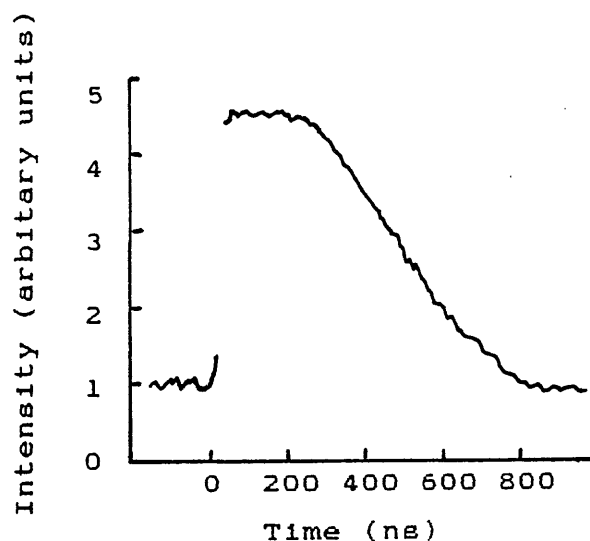
The time at which the experimentally observed fluorescence plateau decayed by 10% was measured to be 270 ± 50 ns (averaged over 10 shots). The CR code (incorporating the time history of the laser pulse) predicts that the 10% reduction will not occur until 700 ns after initiation of the laser pulse. Thus the plateau does appear to decay whilst the Balmer- α transition is still saturated.

ii) 50 us delay. Figure 4.9 shows a typical fluorescence trace and the prediction of the CR code, the plateau heights being normalised for the purpose of comparison. The experimental data was found to decay by 10% 500 ± 70 ns after laser initiation (averaged over 10 shots). Again the CR code predicts that this decay will not occur until approximately 700 ns after initiation.

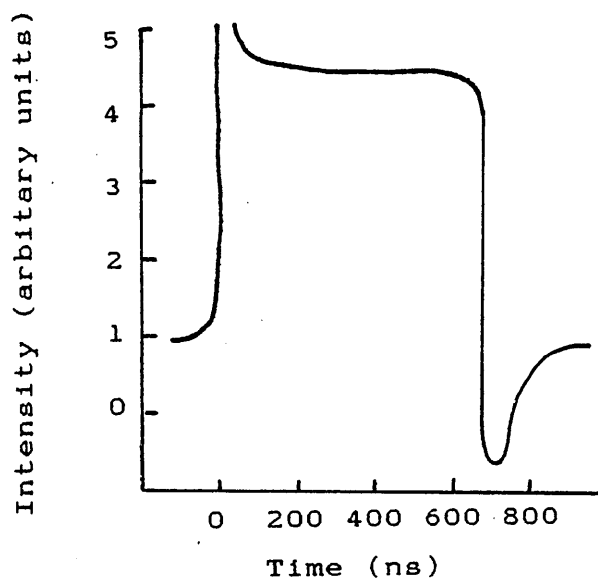
iii) 70 us delay. Figure 4.10 shows a typical fluorescence trace and the prediction of the CR code again normalised for the purpose of comparison. The experimental results display a decay of 10% after 730 ± 100 ns, again averaged over 10 shots. The CR code predicts that the fluorescence signal will not have decayed by 10% until approximately 720 ns after the laser initiation.

The results are presented in Table 4.1 with the photo-detachment fraction (calculated from Equation 4-3) corresponding to the time when the plateau was observed to decay. From this table it can be seen that

Figure 4.8 Balmer-Alpha Fluorescence Trace at 30 us Delay Time.



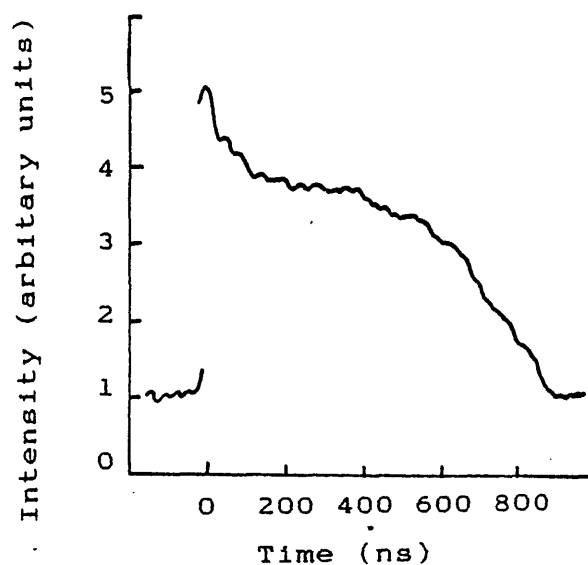
1) Typical Experimental Trace.



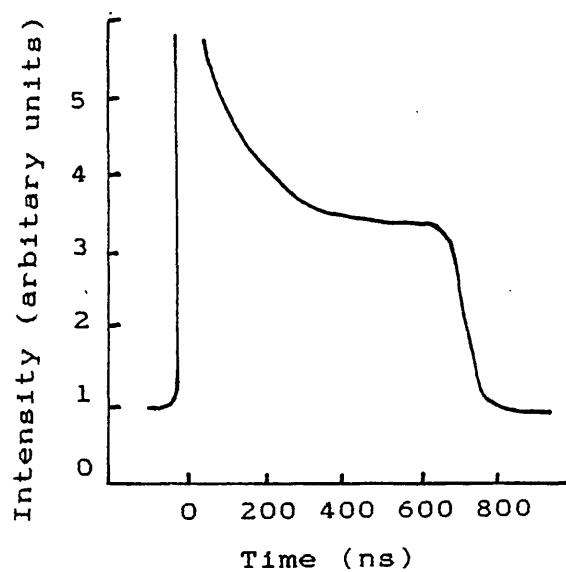
11) CR Code Output.

This is the predicted fluorescence trace from the CR code for the rates of Vriens and Smeets (1980) and a laser time history as shown in Figure 4.7.

Figure 4.9 Balmer-Alpha Fluorescence Trace at 50 us Delay Time.



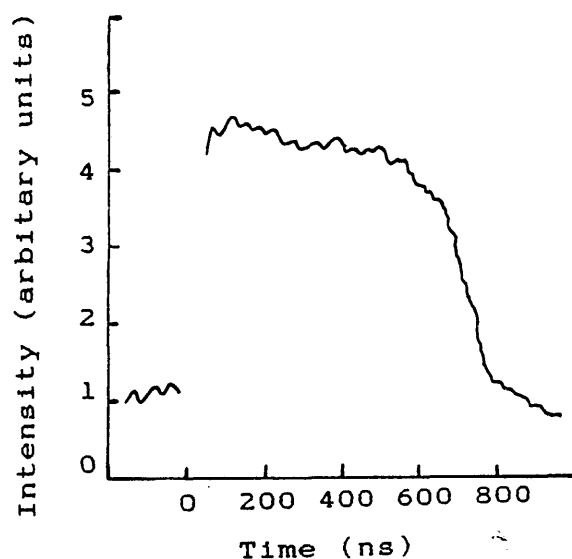
1) Typical Experimental Trace.



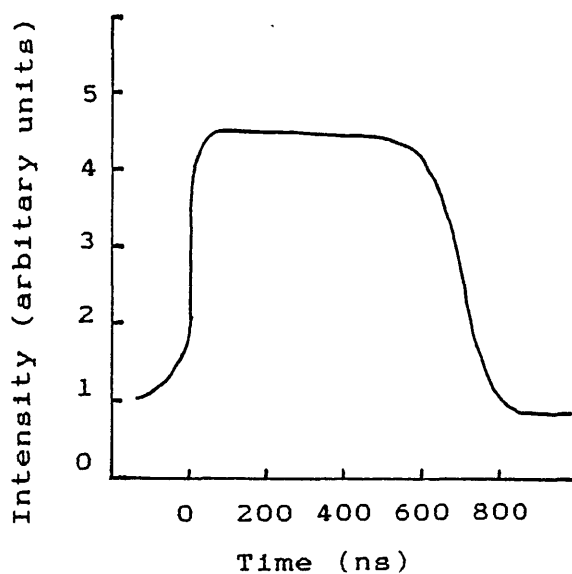
1i) CR Code Output.

This is the predicted fluorescence trace from the CR code for the rates of Vriens and Smeets (1980) and a laser time history as shown in Figure 4.7.

Figure 4.10 Balmer-Alpha Fluorescence Trace at 70 us Delay Time.



1) Typical Experimental Trace.



11) CR Code Output.

This is the predicted fluorescence trace from the Cr code for the rates of Vriens and Smeets (1980) and a laser time history as shown in Figure 4.7.

TABLE 4.1: Time After Laser Initialisation that Plateau Enhancement Decays.

Delay Time (μs)	n_e ($\times 10^{14} \text{ cm}^{-3}$)	$t_{\text{predicted}}$ (ns)	t_{observed} (ns)	$\Delta n_-/n_-$
30	10.70	700	270 ± 50	$0.15 \rightarrow 0.60$
50	2.02	700	500 ± 70	$0.83 \rightarrow 0.92$
70	0.38	720	700 ± 100	$0.95 \rightarrow 1.00$

This table shows $t_{\text{predicted}}$, the time at which the predicted fluorescence behaviour (from the CR code for the rates of Vriens and Smeets, 1980) falls by 10% of its equilibrium value and t_{observed} , the measured time at which the plateau begins to decay from its initial value. Also shown is the photodetachment fraction of the H^- ions ($\Delta n_-/n_-$) corresponding to t_{observed} .

$t_{\text{observed}} < t_{\text{predicted}}$ except for the 70 ns delay time, suggesting that the laser is indeed depressing the plateau enhancement.

4.4 Photodetachment of H⁻ Ions By A Ruby Laser Beam.

The aim of this set of experiments was to irradiate the plasma with a ruby pulse to deplete the H⁻ population and to study the recombination time of both the H- α and H- β emission and the absolute depletion of the emission at these wavelengths.

4.4.1 Overview of the Experimental System.

1) The Laser. The laser used in this experiment was a Korad K1000 series ruby laser, an oscillator/ amplifier system with an intra cavity KDP Pockel cell to provide Q-switching. The pulse had an approximate FWHM of 30 ns providing typically 1.5-2 J of energy at the input window of the pinch. Allowing for 8% loss due to the entrance window this would give 46-61 MW of power at the pinch centre.

By focussing the beam through a variable aperture and measuring the energy as a function of the aperture diameter, a spatial beam profile was obtained. The half angle half power divergence of the laser was measured to be 1.3 ± 0.2 mrad and was found to agree with the manufacturer's quote. The half angle divergence at the 90% power point was measured to be 4.8 ± 1.0 mrad.

ii) Experimental Arrangement. Figure 4.11 is a schematic representation of the experimental arrangement used. The laser beam was fired through the centre of the pinch, the beam being dumped onto a slab of OB10 blue glass at Brewsters angle, with another slab opposite to absorb any reflections. A photodiode (risetime approximately 1 ns) to monitor the laser was placed behind the

blue glass dump with a few tissues placed over it to diffuse the light.

The emission was observed perpendicular to the ruby beam across the diameter of the z-pinch. The light was focussed into the 1 m Hilger Monospek monochromator in the same way as the fluorescence signal in the high power fluorescence experiment described in Section 4.2.1 (i±i) The same C31034A photomultiplier was used (risetime 1 ns at -1400 V) and the emission and laser monitor displayed on two Tektronix 7623A storage oscilloscopes, having quoted bandwidths of 100 MHz.

The ported pinch section described in the previous experiment was originally used, so that a Rayleigh horn could be placed opposite the viewing port to reduce stray light. The emission at both the Balmer-alpha and Balmer-beta wavelengths was observed with this particular vessel for an unrestricted ruby beam at delay times of 30, 40, 50 and 60 us into the recombination phase. The ported vessel was then replaced by the original unported one to eliminate any end zone effect (due to cold plasma in the side arms) on the results. Observations of the Balmer- α emission were taken for delay times of 20 and 30 us into the recombination phase both with and without a 0.5 cm aperture in front of the entrance port.

iii) Timing. The timing sequence began with a pulse from the ruby laser (which was synchronous with the flow of current into the ruby flashlamps) being fed into a Chelsea Instruments delay unit in an electrically screened box. A delay of 1.1 ms was generated by the delay unit to allow a build up of the population inversion in the rods and then a TTL pulse was sent to fire the z-pinch. The pinch photodiode monitor then produced a TTL pulse which was also fed into the delay unit. A delay of 20-60 us could then be generated for the oscilloscope and laser pulse trigger so that the desired time in the plasma recombination phase could be viewed.

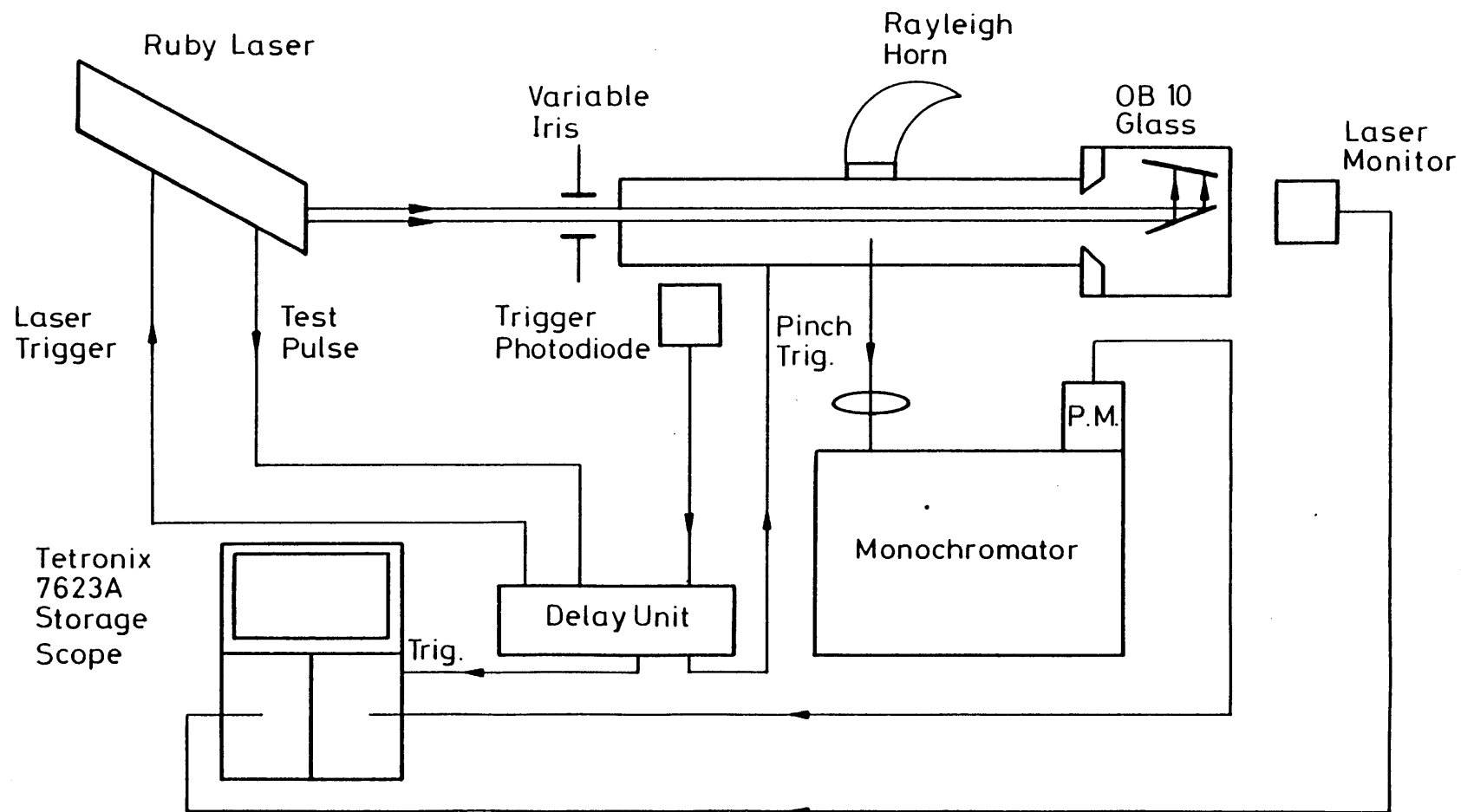


Figure 4.11 Experimental Arrangement for Photodetachment of H^- Ions by Ruby Laser.

iv) The Geometric Factor. The geometric factor g is defined as the ratio of the ruby beam diameter at the plasma centre to the pinch diameter. Figure 4.12 shows a tracing of a typical emission trace observed. The ratio B/A from the diagram should be equal to the geometric factor g provided that no repopulating rate exists for the $n=3$ or 4 level comparable to the ruby photoionization/ H^- photodetachment rate.

For the ported z-pinch section, there is evidence that cold plasma exists in the ports (discussed in Section 4.2.1.3). The extent to which this extra plasma could affect the absolute depletion observed cannot be quantified, but the effective pinch diameter of 4.7 cm would be increased (by an estimated 2 ± 1 cm). For the half power divergence a ruby beam diameter of 1.73 ± 0.03 cm was calculated for the centre of the pinch, giving a geometric factor of 0.26 ± 0.04 . For the unported section, the pinch diameter was measured to be 4.70 ± 0.05 cm. For the unrestricted ruby beam, the half power divergence defined a geometric factor of 0.37 ± 0.05 . The 0.5 cm diameter aperture placed 50 cm in front of the plasma volume viewed, defined the half angle of divergence to be 3.80 ± 0.38 mrad. The geometric factor associated with the aperture was calculated to be 0.19 ± 0.04 from this.

4.4.2 Experimental Results.

1) Balmer-Alpha Experimental Results. Plates IV v) and vi) show typical traces obtained for the conditions studied plus a typical ruby laser monitor trace. The emission profiles shown on Plate IV v) are for the unported z-pinch section and the profiles shown in Plate IV vi) are for the z-pinch section with side ports. Note that the magnitude of the depletion appears to vary with the conditions, becoming larger with increasing delay time (decreasing density and temperature). The time the $H-\alpha$ signal takes to return to its initial level also

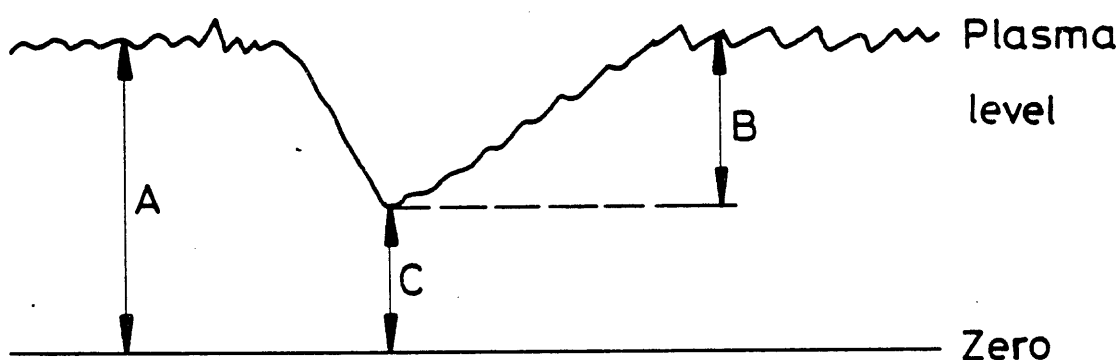


Figure 4.12 Typical Emission Trace Showing Depletion by Ruby Laser.

A is the plasma background level, C is the depletion level (due to photoionization). The ratio B/A is the dip ratio and should be equal to the geometric factor g if all the upper level is removed by the ruby laser.

PLATE IV v)

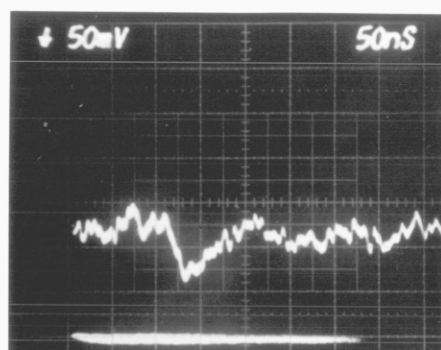
Balmer-Alpha Emission During Ruby Laser Irradiation

Unrestricted laser beam

Delay time = 20.00 μ s

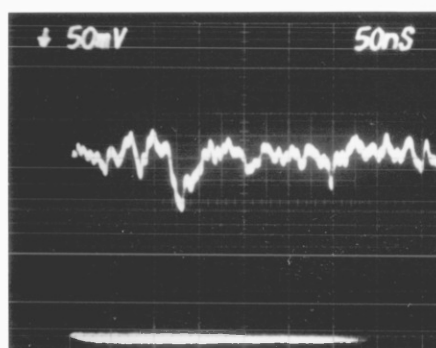
Electron density = $24.60 \cdot 10^{14} \text{ cm}^{-3}$

Temperature = 1.03 eV



5 mm restricted laser beam

Delay time = 20.00 μ s

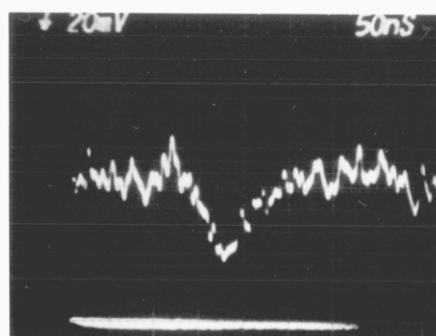


Unrestricted laser beam

Delay time = 30.00 μ s

Electron density = $10.70 \cdot 10^{14} \text{ cm}^{-3}$

Temperature = 0.74 eV



5 mm restricted laser beam

Delay time = 30.00 μ s

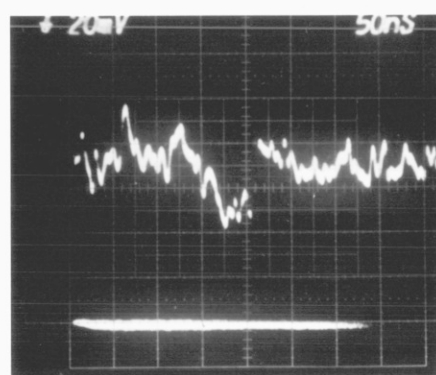
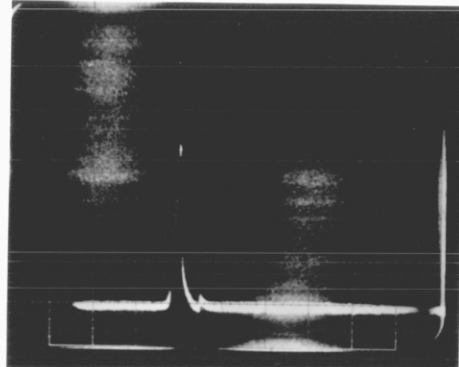


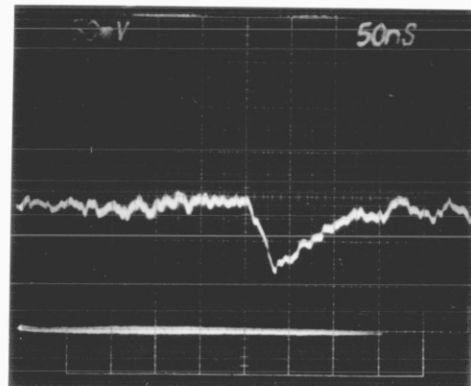
PLATE IV vi)

Balmer-Alpha Emission During Ruby Laser Irradiation
(Unrestricted Laser Beam).

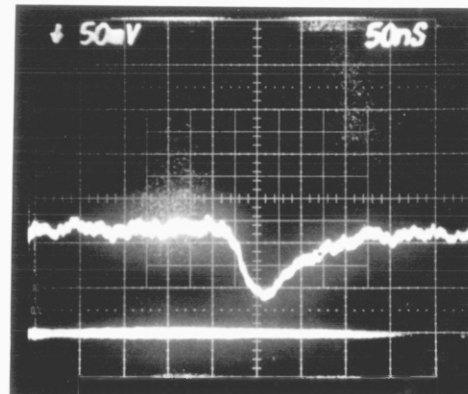
Laser monitor
1 V/div 100 ns/div



Delay time = 40.00 μ s
Electron density = $4.60 \cdot 10^{14} \text{ cm}^{-3}$
Temperature = 0.53 eV



Delay time = 50.00 μ s
Electron density = $2.02 \cdot 10^{14} \text{ cm}^{-3}$
Temperature = 0.38 eV



Delay time = 60.00 μ s
Electron density = $0.90 \cdot 10^{14} \text{ cm}^{-3}$
Temperature = 0.27 eV

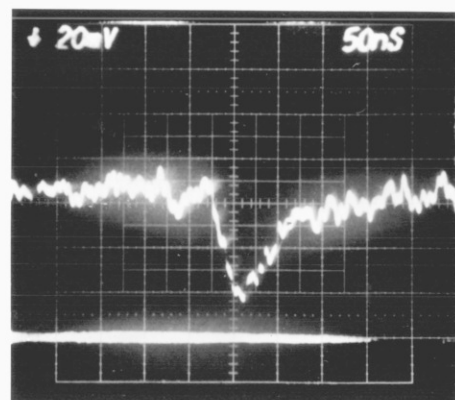
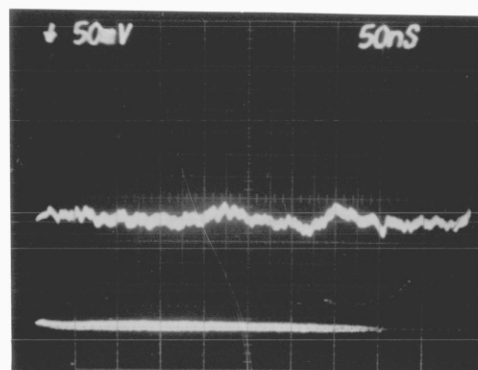


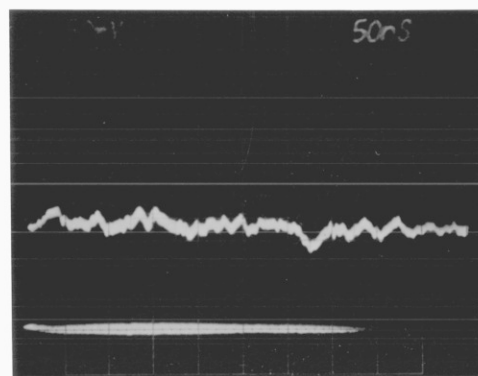
PLATE IV vii)

Balmer-Beta Emission During Ruby Laser Irradiation
(Unrestricted Laser Beam).

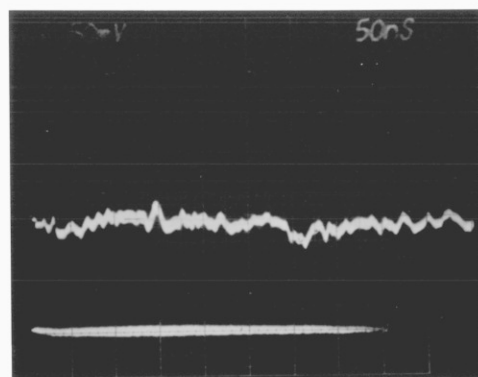
Delay time = 30.00 μ s
Electron density = $10.70 \cdot 10^{14} \text{ cm}^{-3}$
Temperature = 0.74 eV



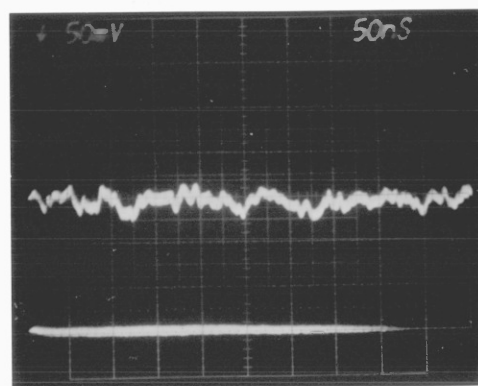
Delay time = 40.00 μ s
Electron density = $4.60 \cdot 10^{14} \text{ cm}^{-3}$
Temperature = 0.53 eV



Delay time = 50.00 μ s
Electron density = $2.02 \cdot 10^{14} \text{ cm}^{-3}$
Temperature = 0.38 eV



Delay time = 60.00 μ s
Electron density = $0.90 \cdot 10^{14} \text{ cm}^{-3}$
Temperature = 0.27 eV



appears to increase with decreasing density and temperature. The data is discussed in more detail in the next section.

ii) Balmer-Beta Experimental Results. The monochromator was tuned to the H- β (4861 Å) transition to observe the effect of the ruby pulse on the n=4 level. This was to eliminate any optical depth effects, H- β being an optically thin transition at all the conditions investigated. All results at this wavelength were taken with an unrestricted ruby beam in the ported z-pinch section.

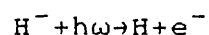
Plate IV vii) shows typical traces obtained for delay times of 30, 40, 50 and 60 us after peak current flow. Note that the actual depletion of the level observed is much smaller than for the H- α data under the same pinch conditions, and that the H- β emission returns to its background level in a much faster time.

4.5 Discussion of Experimental Results.

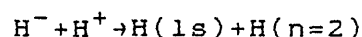
4.5.1 Relevant Atomic Processes. As discussed in relation to the H- α laser experiment in Section 4.2.2.4, continuum absorption was considered to be negligible. However, for a discussion of the results from this experiment, the rates associated with the relevant processes must be considered.

1) H- Population. The main processes associated with the H⁻ ions in this plasma are:

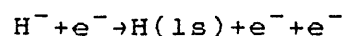
a) Photodetachment of the H⁻ ions via the ruby laser.



b) Depletion of H⁻ ions via neutralization feeding the n=2 and 3 levels.



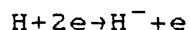
c) Collisional detachment by electrons.



d) Collisional recombination feeding the H^- ion population.



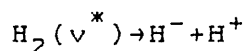
e) Collisional attachment feeding the H^- ion population.



During ruby laser irradiation process a) is the dominant process. The cross section for this process at the ruby laser wavelength is $4 \cdot 10^{-17} \text{ cm}^2$ (Takayanagi and Suzuki, 1978). Given a typical power of 50 MWcm^{-2} this gives a photodetachment rate of $7 \cdot 10^9 \text{ s}^{-1}$ per H^- ion. The maximum rate per H^- ion for processes b) is $7.4 \cdot 10^6 \text{ s}^{-1}$ into the $n=2$ level and $7.4 \cdot 10^7 \text{ s}^{-1}$ into the $n=3$ level for the 20 μs delay time (Bates and Lewis, 1955). The maximum rate per H^- ion for process c) is $9.8 \cdot 10^6 \text{ s}^{-1}$ for the 20 μs delay time (Takayanagi and Suzuki, 1978). Process d) is the largest repopulating rate for the H^- ions under the conditions investigated, having a maximum rate per ground state atom of 8.6 s^{-1} and process e) has a maximum rate of approximately 4.8 s^{-1} per ground state atom. Thus the ruby laser photodetachment dominates the processes associated with the H^- ions.

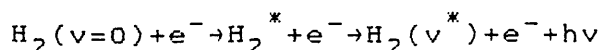
From Equation (4-3), the photodetachment fraction $\Delta n_-/n_-$ is unity and all the H^- ions are therefore considered to be removed from the plasma by the ruby laser pulse.

Note that the production of H^- ions in multipole "bucket" sources such as exist at UKAEA Culham Laboratory (Hiskes et al, 1983) is thought to be dominated by the dissociative attachment of vibrationally excited molecules with vibrational quantum number v between 6 and 12 (Hiskes and Karo, 1984):



At 1 eV, the rate coefficient per $H_2(v=8)$ molecule is approximately 10^7 s^{-1} (Wadehra, 1979). To compete with the main H^- populating processes at 1 eV, the $H_2(v^*)$

density would need to be of the same density as the LTE H^- density ($2 \cdot 10^{10} \text{ cm}^{-3}$, from Table 2.6). The main process feeding $H_2(v^*)$:

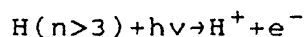


is negligible under the plasma conditions studied here since the molecular hydrogen has to collide with electrons of energy greater than 20 eV (Hiskes and Karo, 1984).

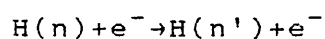
After termination of the ruby laser pulse, the depleted H^- ion population will recombine to its initial, pre-laser, level in a time set only by the depopulating rates (processes (b) and (c)). Table 4.2 shows the time required for the H^- population to return to e^{-1} of its initial value, for the plasma conditions studied and using the rates of Bates and Lewis (1955). These recombination times vary between 12 and 265 ns.

11) The Atomic Hydrogen Population. The ruby laser also photoionizes the $n \geq 3$ levels in atomic hydrogen, and so the main rates associated with these levels must also be considered. They are:

a) Photoionization of the $n \geq 3$ levels in atomic hydrogen.



b) Electron collisional excitation and de-excitation between the relevant levels.



The cross section for the photoionization of the $n=3$ level at 6943 Å is $1.4 \cdot 10^{-17} \text{ cm}^2$ from Takayanagi and Suzuki (1978) giving a photoionization rate of $2.8 \cdot 10^9 \text{ s}^{-1}$ per $n=3$ atom. The electron collision rates out of the $n=3$ level over the conditions considered vary from $1.5 \cdot 10^9$ to $7.7 \cdot 10^6 \text{ s}^{-1}$ per $n=3$ atom (Vriens and Smeets, 1980). The photoionization rate out of level 3 is at least an order of magnitude higher than the main rate out of level

TABLE 4.2: Recombination Time for the H^- Population After Complete Depletion by Ruby Pulse Irradiation.

Delay Time (us)	n_e (10^{14} cm^{-3})	Recombination Time (ns)
20	24.6	12.3
30	10.7	27.5
40	4.6	55.7
50	2.0	125.0
60	0.9	264.6

TABLE 4.3: The Repopulation Time for the $n=3$ Level After Complete Depletion by Ruby Pulse Irradiation.

Delay Time (us)	n_e (10^{14} cm^{-3})	Repopulation Time from collisions (ns)
20	24.6	0.9
30	10.7	2.6
40	4.6	7.0
50	2.0	26.7
60	0.9	89.4

Note that the collision rates from Vriens and Smeets have been used to calculate the repopulation time of the $n=3$ level in this case.

3 for the delay times of 30, 40, 50 and 60 us. At a delay time of 20 us, since there is much evidence to suggest that the collision rates out of the $n=3$ level are approximately a factor of 10 smaller than those given by Johnson (see Section 2.2.1 (v)), then the ruby laser photoionization ought to be the largest rate out of $n=3$ for these times also. The ruby laser was then assumed to deplete the entire $n=3$ level population for all the times considered.

For the higher levels, the photoionization cross section varies as n^{-5} (the Kramers formula) and so the photoionization rate per $n=4$ atom is approximately $5.5 \cdot 10^8 \text{ s}^{-1}$ and per $n=5$ atom is $1.8 \cdot 10^8 \text{ s}^{-1}$. The electron collision rates out of the $n=4$ level vary between $8.8 \cdot 10^9 \rightarrow 5.7 \cdot 10^7 \text{ s}^{-1}$ (Vriens and Smeets, 1980) and for the $n=5$ level between $2.8 \cdot 10^{10} \rightarrow 2.8 \cdot 10^8 \text{ s}^{-1}$ (Vriens and Smeets) over the plasma conditions investigated. The $n=4$ level population may not be completely depleted by the ruby laser (particularly for the earlier delay times), depending on the actual values of the collision rates out of that level. For the $n \geq 5$ levels, depletion will hardly occur at all, the collision rates being much larger than the ruby photoionization rate.

After termination of the laser pulse, the $n=3$ and 4 level populations will return to their initial (pre-laser) levels in a time given by the dominant rates out of these levels. The $n=3$ level will be repopulated in a time given by the inverse of the $n=3 \rightarrow 2$ and $3 \rightarrow 4$ collision rates. Table 4.3 shows the repopulation time (assuming a constant $n=2$ density) for the $n=3$ level to return to e^{-1} of its initial value. From this table it can be seen that the repopulation times ($1 \rightarrow 89 \text{ ns}$) of the $n=3$ level are much shorter than those of the H^- population (shown in Table 4.2). Note that since the oscilloscope was operated at a timescale of 50 ns/cm the minimum time interval that could be confidently measured was approximately 10 ns . Repopulation times of less than this could therefore not

be accurately measured.

As previously mentioned, not all of the $n=4$ level population will be photoionized due to competition from the electron collision rates. To investigate the depletion of the $n=4$ level for the delay times investigated, consider the main rates into and out of the $n=4$ level (the $n=4 \rightarrow 5$ and the $5 \rightarrow 4$ rates). The $n=5$ level population is assumed to remain constant in this case. The rate of change of the $n=4$ level may be approximated to:

$$\frac{dn_4}{dt} = K_{54}n_5 - K_{45}n_4 \quad (4-4)$$

where K_{45} and K_{54} are the collisional excitation/de-excitation rates s^{-1} per atom. Thus the $n=4$ population for a given time t after turning on the laser pulse (with a photoionization rate R out of the $n=4$ level) is:

$$n_4^t = \frac{K_{54}n_5}{(K_{45} + R)} \quad (4-5)$$

Table 4.4 shows the approximate dip ratio B/A (defined in Figure 4.12) expected from the above analysis for the Balmer- β transition. Note that the $n=5$ population is obtained from the CR code under steady state conditions and the dip ratios are calculated from both the rates of Johnson and Vriens and Smeets.

The $n=4$ level will repopulate in a time given by the rate out of that level (essentially the $n=4 \rightarrow 5$ collision rate). Table 4.5 shows the repopulation times of the $n=4$ level for the range of plasma conditions studied. These are much faster than the recombination times of the H^- ions as shown in Table 4.2 and also the $n=3$ repopulation times, shown in Table 4.3.

In summary, careful consideration of the relevant atomic processes and associated rates shows that all the H^- ions present would be removed from the plasma due to

TABLE 4.4: The Dip Ratio B/A for the n=4 Level During Ruby Irradiation.

Delay Time (us)	n_e (10^{14} cm^{-3})	Dip Ratio B/A		Dip Ratio Geometrically Corrected	
		Johnson	Vriens	Johnson	Vriens
20	24.6	0.16	0.20	0.06 $\pm$ 0.02	0.08 $\pm$ 0.02
30	10.7	0.16	0.26	0.06 $\pm$ 0.02	0.10 $\pm$ 0.03
40	4.6	0.16	0.39	0.06 $\pm$ 0.02	0.15 $\pm$ 0.04
50	2.0	0.37	0.88	0.14 $\pm$ 0.04	0.34 $\pm$ 0.09
60	0.9	0.64	0.95	0.25 $\pm$ 0.06	0.37 $\pm$ 0.10

TABLE 4.5: The Repopulation Time for the n=4 Level Population After Termination of the Ruby Laser.

Delay Time (us)	n_e (10^{14} cm^{-3})	Repopulation Time from collisions (ns)
20	24.6	0.2
30	10.7	0.5
40	4.6	1.2
50	2.0	4.0
60	0.9	12.8

Note that the collision rates from Vriens and Smeets (1980) have been used to calculate the repopulation time of the n=4 level in this case.

photodetachment by the ruby laser, and would not recombine to within e^{-1} of the pre-laser level until 12→265 ns had elapsed after termination of the ruby laser pulse. The laser was also assumed to photoionize all of the $n=3$ level population in the atomic hydrogen, and some of the $n=4$ level population, the main collision rates out of this level being comparable to or larger than the photoionization rate out. The repopulation times for both the $n=3$ and 4 level populations were shown to be much faster than those for the H^- ion population.

4.5.2 Balmer-Alpha Emission Results.

1) The Observed Dip. Figure 4.13 shows the magnitude of the measured depletion/background ratio B/A (defined in Figure 4.12) as a function of electron density for the unrestricted ruby beam and the ported pinch section. Also shown in Figure 4.13 are the theoretical values as given by the geometric factor for the measured half power divergence of the ruby beam as discussed in Section 4.4.1 (iv). Each data point represents an average of at least 10 shots and the error bars on the data are the standard deviations about the mean. No correction has been made for the finite optical depth at $H-\alpha$. The data points for the 40, 50 and 60 μs delay times ($n_e = 0.9 \text{ } 10^{14} \rightarrow 4.6 \text{ } 10^{14} \text{ cm}^{-3}$) consistently lie above the range given by the geometric factor, a larger depletion being observed. The data from the 30 μs delay time ($n_e = 10.7 \text{ } 10^{14} \text{ cm}^{-3}$) lies in the predicted range of values but the 20 μs delay data ($n_e = 24.6 \text{ } 10^{14} \text{ cm}^{-3}$) lies above the predicted values.

Figure 4.14 shows the depletion/background ratio B/A measured as a function of the electron density in the unported z-pinch for the unrestricted and 0.5 cm aperture restricted ruby beam. For the unrestricted laser beam, the data for the 30 μs delay ($n_e = 10.7 \text{ } 10^{14} \text{ cm}^{-3}$) lies above the predicted depletion and the 20 μs delay ($n_e = 24.6 \text{ } 10^{14} \text{ cm}^{-3}$) within the theoretical range of values. The depletion in the $H-\alpha$ emission observed with

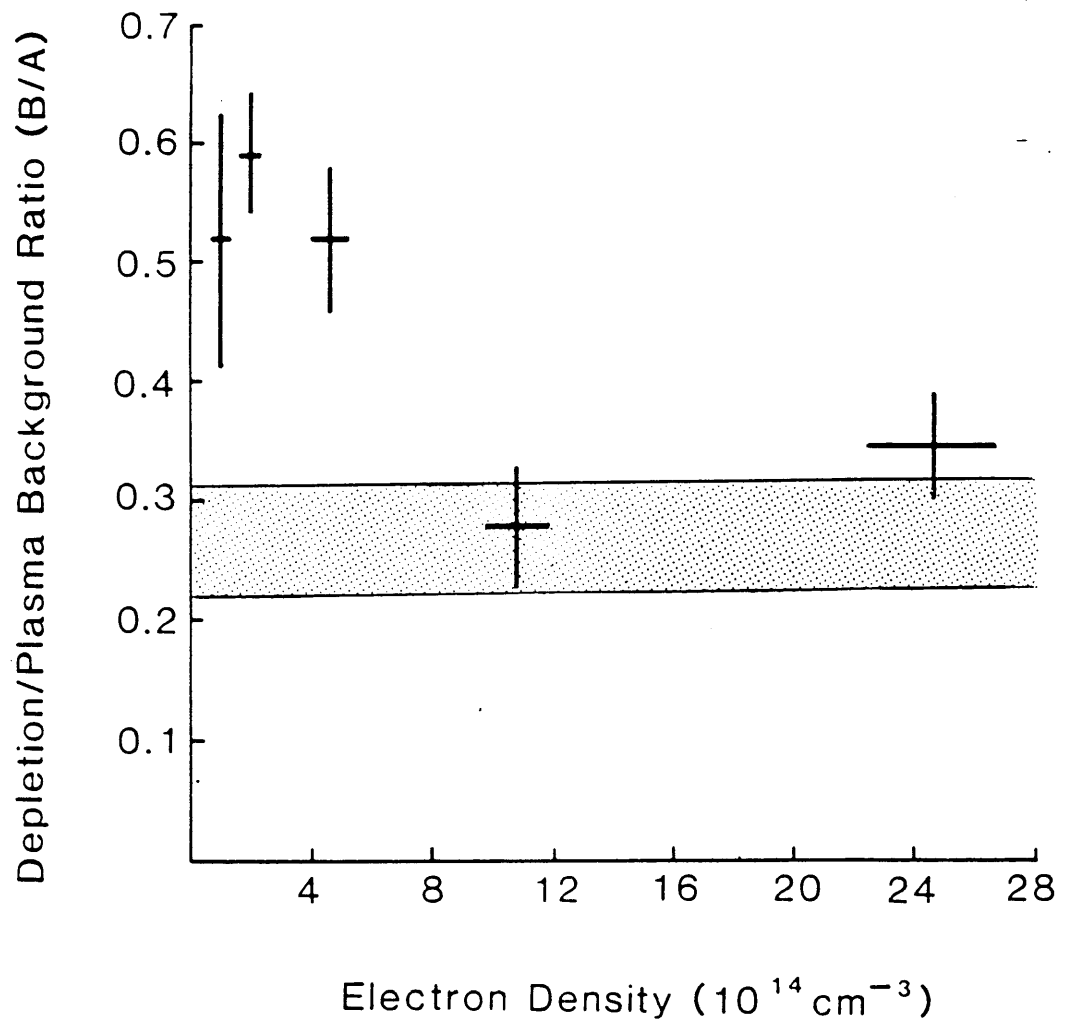


Figure 4.13 The Depletion of the Observed Balmer-Alpha Emission as a Function of Electron Density (Ported Pinch Section).

This shows the measured depletion/background ratio B/A (defined in Figure 4.12) for the unrestricted ruby beam and the ported pinch section. The continuous region represents the predicted dip calculated from the geometric ratio set by the half power divergence of the ruby laser for the full aperture.

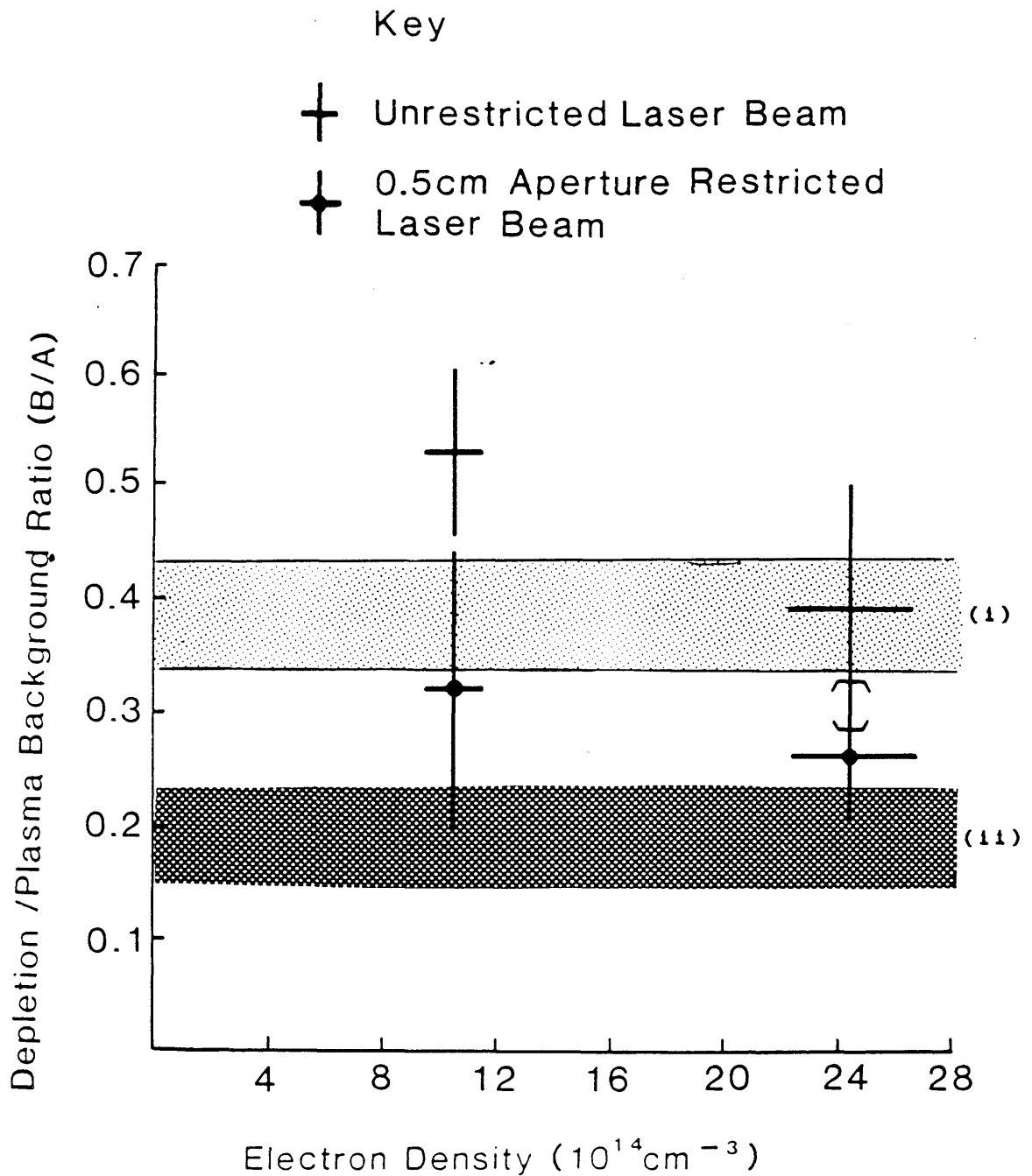


Figure 4.14 The Depletion of the Observed Balmer-Alpha Emission as a Function of Electron Density (Unported Pinch Section).

This shows the measured depletion/background ratio B/A (defined in Figure 4.12) for the unrestricted and 0.5 cm aperture restricted ruby beam and the unported pinch section. The continuous regions represent the predicted dip calculated from the geometric ratio set by the half power divergence of the ruby laser for:

- i) the full aperture
- ii) the 0.5 cm aperture in front of the beam.

the well defined beam lies above the predicted value for both the 30 and 20 us delay times although the lower part of the error bars do overlap the predicted value. Note that no correction has been made for the finite optical depth at the H- α wavelength. The effect of this would be to reduce the background level and so decrease the predicted depletion in the emission. This would increase the discrepancy observed for some of the data points in Figures 4.13 and 4.14.

In the cases where the observed depletion is more than that predicted, it is necessary to consider whether these results could be due to inhomogeneities in the ruby laser beam, or by diffusion of atoms from unpumped regions into the irradiated volume.

As regards the first option, the power divergence of the laser was actually measured (Section 4.4.1 (1)). When the beam size was well defined by an aperture, the depletion of the H- α emission was still generally observed to be higher than predicted.

For the second point, consider a thermal atom (this would give an overestimate of the diffusion rate). The maximum distance it would travel in the time duration of the ruby pulse (30 ns) would be 415 μm for the 20 us delay time. This is not a large enough increase in the size of the beam to account for the discrepancies observed. For the later times, the atoms would have travelled a smaller distance.

ii) Repopulation Time. The time for the H- α emission to return to e^{-1} of its initial level after termination of the laser pulse was measured. Table 4.6 shows the measured repopulation times for the plasma conditions considered, and range from 27.5 ± 9.8 to 80.6 ± 19.0 ns.

Comparison of the predicted time for the H $^-$ ions to recombine (Table 4.2) with the experimentally observed

TABLE 4.6: The Observed Time for the Balmer-Alpha
Emission to Return to 1/e of its Initial Level After
Termination of the Ruby Pulse.

Delay Time (us)	n_e (10^{14} cm^{-3})	n=3 Repopulation Time (ns)
20	24.6	27.5 $\pm$ 9.8
30	10.7	37.9 $\pm$ 6.0
40	4.6	59.2 $\pm$ 7.0
50	2.0	78.0 $\pm$ 5.0
60	0.9	80.6 $\pm$ 19.0

repopulation time (Table 4.2) shows that the H- α emission returns to its initial level slower than the H⁻ ions recombine for the 20 and 30 us delay data, in roughly the same time for the 40 us data and faster for the 50 and 60 us delay time data. Thus the n=3 density may be fed in part by the H⁻ ion population at the earlier times but not at the later ones.

Comparison of the measured repopulation time and the predicted repopulation time for collisions from the n=2 and 4 level (Table 4.3) shows that the measured repopulation times are slower than predicted. Thus if the repopulation times observed are due to collisions from the n=2 level, the total rate out of n=3 (which is dominated by the n=3 $\rightarrow$ 4 rate) would have to be factors of between 3 to 6 slower than even that given by Vriens and Smeets (except for the 60 us delay time when the observed and predicted times agree to within experimental error).

iii) Black Body Radiation. No account has yet been made of the black-body radiation field at the Balmer- α wavelength. From Thorne (1974), the thermal n=3 $\rightarrow$ 2 rate is given by:

$$R_{32} = B_{32}p(\nu_{32}) \quad (4-6)$$

where B_{32} is the Einstein stimulated emission coefficient and $p(\nu_{32})$ is the thermal radiation field density at Balmer- α . Now

$$B_{32} = A_{32}(c^3/8\pi h\nu^3) \quad (4-7)$$

where A_{32} is the spontaneous 3 $\rightarrow$ 2 decay rate (4.4×10^7 s⁻¹), and all other symbols have their usual meanings. The radiation density is given by the black body function:

$$p(\nu_{32}) = (8\pi h\nu^3/c^3)(e^{h\nu/kT} - 1)^{-1} \quad (4-8)$$

Substituting equations (4-7) and (4-8) into (4-6) gives the following expression for the time taken for the Balmer- α black body field to repopulate the depleted $n=3$ level to within e^{-1} of its initial level:

$$R^{-1} = (e^{(h\nu/kT)} - 1) / A_{32} \quad (4-9)$$

Table 4.7 shows this time calculated for the plasma conditions studied here. (Note that the corresponding time for the Lyman- β black body field would be 21 us (at the 20 us delay time), assuming this line to be optically thin.)

Delay Time (us)	n_e (10^{14} cm^{-3})	R^{-1} (us)
20	24.6	0.12
30	10.7	0.19
40	4.6	0.78
50	2.0	33.0
60	0.9	248.0

TABLE 4.7: The Time Taken for the Black-Body Balmer-Alpha Field to Repopulate the Depleted $n=3$ Level to $1/e$ of its Initial Value.

The times given in Table 4.7 are longer than the measured repopulation times by factors of at least 5, and so the thermal Balmer- α field is not responsible for the observed $n=3$ repopulation time.

4.5.3 Balmer-Beta Emission Results.

1) The Observed Dip. Figure 4.15 shows the depletion to background ratio B/A measured as a function of the electron density for the ported z-pinch section and the unrestricted laser beam. The errors shown for the single shot point ($n_e = 0.9 \cdot 10^{14} \text{ cm}^{-3}$) are based on the variation in the plasma background for that shot. Also shown are the geometric factors defined by the half power divergence as discussed in Section 4.4.1 (iv) and the calculated depletion ratios from Table 4.5 employing the rates of both Johnson and Vriens and Smeets. Note that these calculations are only an approximation. Error bars have been added for a 10% uncertainty in the diagnosed electron temperature (Section 4.1.4). Note that the H- β transition is optically thin and therefore no correction to the data is required.

The two data points at densities 4.6 and $10.7 \cdot 10^{14} \text{ cm}^{-3}$ are better described by the calculation using the lower rates of Vriens and Smeets, and the two data points at the lower densities of 2 and $0.9 \cdot 10^{14} \text{ cm}^{-3}$ are better described by the calculation using the rates of Johnson. This result is interesting because it is further proof that for the higher densities ($n_e \geq 4.6 \cdot 10^{14} \text{ cm}^{-3}$), the actual inter level collision rates are lower than those given by Johnson (see Chapter 2). (This result refers to the rates between the $n=4$ and 5 levels.)

1i) Repopulation Time. The measured time for the Balmer- β emission to return to within e^{-1} of its initial level is shown in Table 4.8 for the conditions investigated. The times range from less than 10 to 18.4 ns and are much faster than the recombination times of the H^- ions as in the case of the H- α emission.

The repopulation times shown in Table 4.8 are larger than the times expected from collisions down from the $n=5$ level (Table 4.5). However these times (calculated from

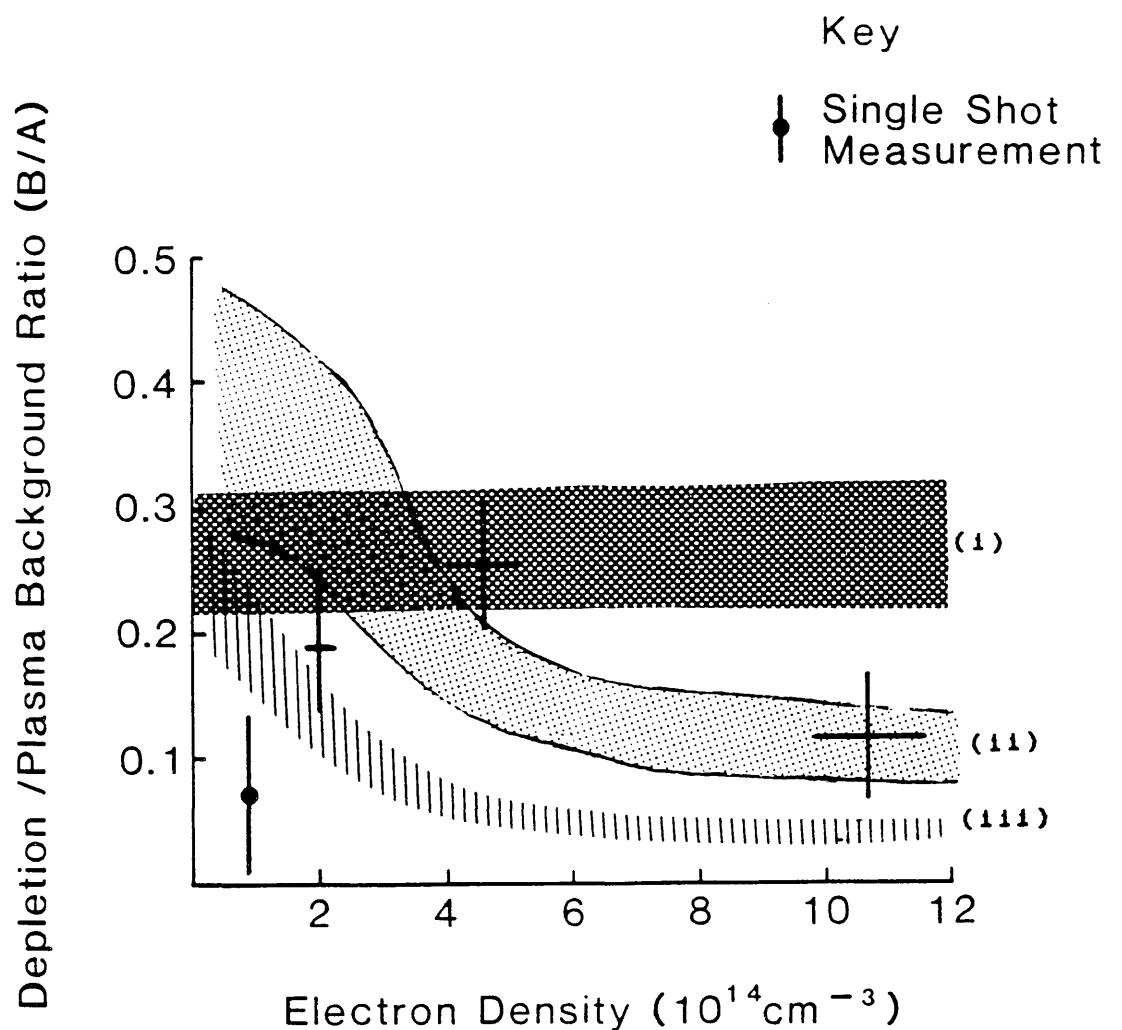


Figure 4.15 The Depletion of the Observed Balmer-Beta Emission as a Function of Electron Density (Ported Pinch Section).

This shows the measured depletion/background ratio B/A (defined in Figure 4.12) for the unrestricted ruby beam and the ported pinch section. The continuous regions represent the predicted dip calculated from:

- i) the geometric factor
- ii) the CR code using Johnson's electron collision rates
- iii) the CR code using Vriens and Smeets' electron collision rates.

TABLE 4.8: The Observed Time for the Balmer-Beta
Emission to Return to 1/e of its Initial Level After
Termination of the Ruby Pulse.

Delay Time (us)	n_e (10^{14} cm^{-3})	Recombination Time (ns)
20	24.6	<10
30	10.7	<10
40	4.6	18.4 ± 6.1
50	2.0	<10
60	0.9	<10

the collision rates) were below the limit of experimental resolution and no further comment can be made about possible discrepancies with theory.

4.6 Conclusions.

The results from the first, high power, fluorescence experiments show the following features:

1) A clear H- α plateau feature, similar to that previously observed by Burgess et al (1980) in the same z-pinch plasma device. The very existence of this feature indicates a basic discrepancy with theory, as discussed in Chapter 2.

ii) The decay of the H- α fluorescence plateau (for the earlier delay times) whilst the laser was still saturating the transition. This could be due to the removal of some species responsible for the artificially high plateau. From work presented in Chapter 2, this species is most likely to be H $^-$ ions. However, it has been shown that mutual neutralization of H $^-$, which populates the n=2 and 3 levels of neutral hydrogen, has little effect upon the H- α fluorescence plateau height (Section 2.3.3). For the premature plateau decay to be due to the removal of a proportion of the H $^-$ population, then either the H $^-$ density is larger than calculated by assuming it to be in LTE with the ground state neutral density (which is true of several excited states of neutral hydrogen), or the mutual neutralization cross section is much larger than calculated.

The results from the second, ruby laser experiment show the following:

i) The observed depletion of the n=3 level by the ruby laser for a well defined beam was larger than expected although the error bars on the experimental data just overlapped the theory.

ii) The $n=3$ level was observed to return to its initial level in a time slower than or approximately equal to the calculated H^- recombination time, and slower than the calculated collision rate for delay times of 20-40 μs . For the 50 μs delay time the measured repopulation time was faster than the H^- ion recombination but slower than the time set by the main collision rates. The measured 60 μs delay time repopulation agreed with the collision rate calculation (using the rates of Vriens and Smeets) but was faster than the H^- recombination. It is therefore unlikely that processes involving H^- ions were responsible for populating the $n=3$ level for these later times. For the earlier times, for electron collision processes to be responsible for the observed repopulation, the $n=3 \rightarrow 4$ and $3 \rightarrow 2$ rates would have to be factors of 3 to 6 times smaller than those given by Vriens and Smeets. This would be in agreement with previous work which indicated that the electron collision rates are actually lower than theoretically predicted (see Chapter 2).

iii) The $n=4$ level was observed to repopulate faster than the calculated H^- recombination, and it may therefore be concluded that $n=4$ is not significantly populated by processes involving this species. Due to the limited time resolution available it was not possible to study the time history of the repopulation of this level further.

Chapter Five

INTRODUCTION TO HIGH TEMPERATURE WORK.

In this chapter, a brief description of the two main toroidal plasma devices used for fusion research (the tokamak and the reverse-field pinch) is given and the reasons for choosing HBTX1A as the device upon which to use the fluorescence scattering technique are discussed. The application of fluorescence scattering as a diagnostic for high temperature fusion-type plasmas is discussed and an overview of previous work in this area is presented.

5.1 Reasons for Neutral Density Measurement on HBTX1A.

It was decided to set up a fluorescence scattering experiment on HBTX1A (High Beta Torus Experiment 1A), a reverse-field pinch at the UKAEA Culham Laboratory. The HBTX1A device was chosen because of its relative flexibility and fast firing rate (approximately 1 shot every 2 minutes being possible). Also, a previous fluorescence scattering experiment had been performed on the same machine by a joint Culham/Imperial College team which produced a surprising result and this required further investigation. The aim of the work on HBTX1A was primarily to obtain a spatially resolved measurement of the neutral density profile, but also to demonstrate that the technique was a reliable and flexible diagnostic for other measurements and could be adapted for routine use on other devices.

The previous experiment on the HBTX 1A device to measure the neutral density at one position in the plasma was performed on HBTX1A by Gohil et al (1982), pumping on the Balmer-alpha transition. This produced an apparently spurious result. Although the ground state density calculated from the fluorescence signal was not unexpected, the actual enhancement observed was too large by a factor of approximately 4 for the given plasma conditions (in comparison with the results from the 20 level collisional-radiative (CR) code as described by Gohil (1982)). Uncertainty in the electron temperature at the spatial position chosen was discounted. The temperature would have had to be as low as 10 eV to account for the actual enhancement observed whereas Thomson scattering measurements indicated 100 eV at the plasma centre. Either the background emission was too low, or the actual fluorescence enhancement was too high, or both.

Gohil (1982) postulated that one reason for this discrepancy could be that the spatial

distribution of the neutrals was not as expected. (An absence of neutrals in the viewing volume would decrease the observed background.) In order to check both this distribution (there being no other neutral measurements at the time) and the signal level, it was decided to first repeat and then extend the work using this technique on HBTX1A.

It should be noted that the calculations by Gohil et al (1983) are in error. The electron density used in their calculations was the peak density for the centre of the plasma ($3 \times 10^{13} \text{ cm}^{-3}$) instead of 10^{13} cm^{-3} corresponding to the relevant position in the plasma (assuming a parabolic distribution of the electron density profile), and "rounding up" errors were also present. The quoted neutral density of $(2.5 \pm 0.4) \times 10^9 \text{ cm}^{-3}$ should in fact be $(7.0 \pm 1.2) \times 10^9 \text{ cm}^{-3}$ when these corrections are made.

The fluorescence scattering measurements presented in this thesis were performed at both the D- α and D- β wavelengths. The reason for moving to the D- β transition was partly to improve the signal to noise ratio (due to the initial population ratios) but mostly in an attempt to obtain a spatial measurement of the electron density, which is discussed in more detail in Chapter 6.4.3.

5.2 Torus Geometry.

Figure 5.1 shows the geometry of a torus defining the major and minor radii. A toroidal field is not sufficient by itself for plasma confinement because there is always some drift of particles in the non uniform field across the minor diameter. Addition of a second (poloidal) field around the minor diameter prevents this.

The two types of magnetically stabilised pinch considered the most likely candidates for a potential fusion reactor are the tokamak and the reverse-field pinch.

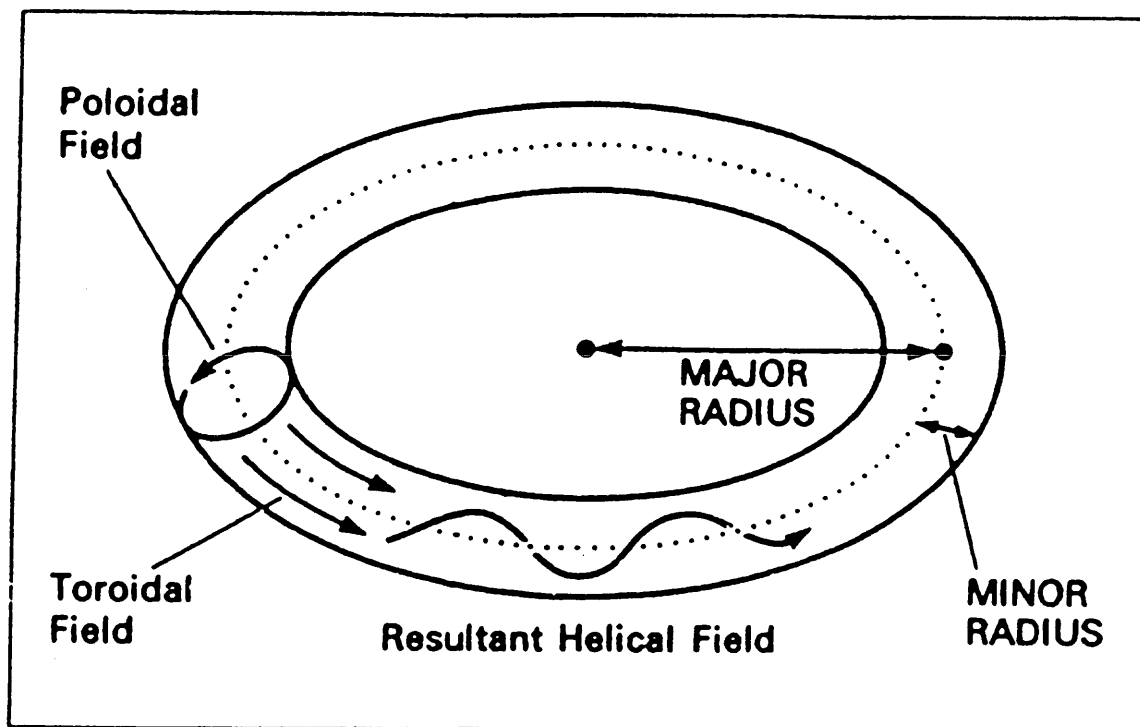


Fig 5.1 *Principal magnetic fields in toroidal geometry.*

(Taken from Culham Laboratory Annual Report, 1983)

TOKAMAK

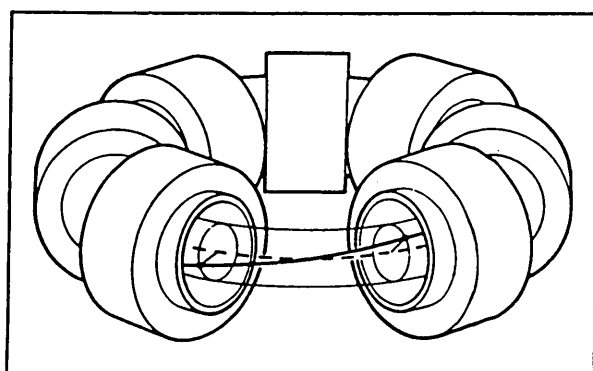


Fig.5.2.1 *The tokamak configuration, showing the gradual twist of the magnetic field lines.*

REVERSED-FIELD PINCH

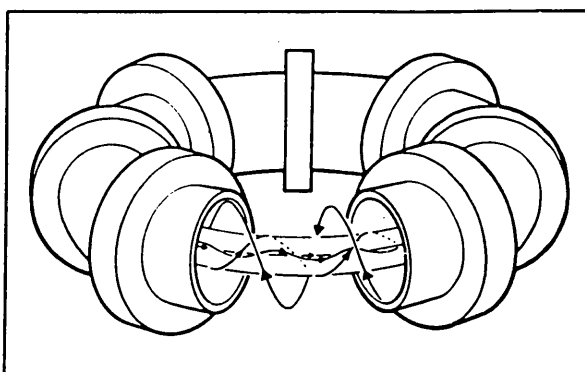


Fig.5.2.2 *The reversed-field pinch configuration, showing (from the centre) the rapid twist of the magnetic field lines including a reversal in the outer region.*

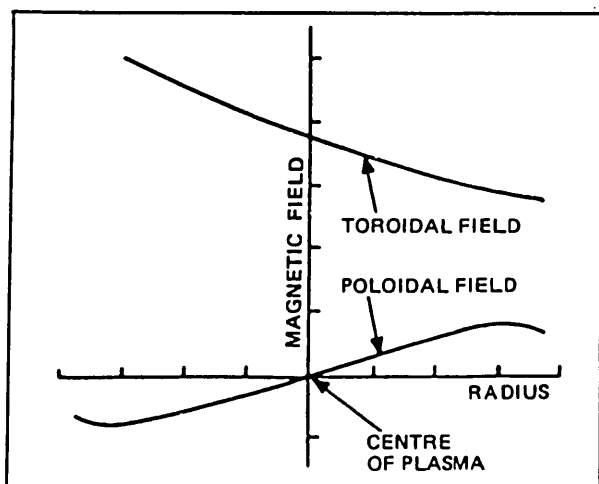


Fig.5.3.1 *The tokamak configuration, showing the variation of the toroidal and poloidal fields as a function of the minor radius. The toroidal field is much stronger than the poloidal field.*

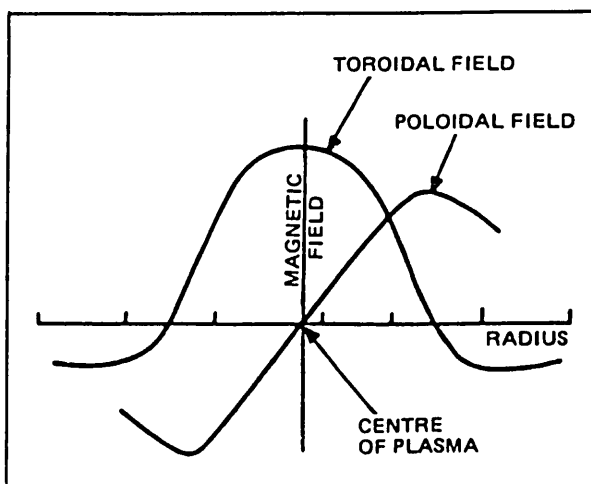


Fig.5.3.2 *The reversed-field pinch configuration, showing the variation of the toroidal field component as a function of the minor radius, showing field reversal in the outer regions. The poloidal field is of the same order of magnitude as the toroidal field.*

characterised by the parameter q , the "safety factor." This is the number of times a field line circles the major axis in circling once round the minor axis. For a magnetic surface of circular cross section, q may be defined as:

$$q = \frac{r B_\phi}{R B_\theta} = \frac{r^2 B_\phi}{2 I R} \quad (5-1)$$

where r and R are the minor and major radii respectively, I is the current in the coils producing the poloidal field B_θ and B_ϕ is the toroidal magnetic field. For the tokamak device, $q > 1$ at all radii and for the reverse-field pinch, $q < 1$ at all radii.

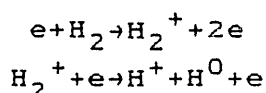
The configurations of the two types of discharge are shown in Figure 5.2. As can be seen, the tokamak configuration gives a gradual twist to the magnetic field lines, and the reverse-field pinch a more rapid twist. The variation of toroidal and poloidal fields as a function of minor radius for both types of discharge is shown in Figure 5.3. For the tokamak system, the toroidal field is much stronger than the poloidal field but for the reverse-field pinch system they are both of the same order of magnitude. Note also that in the latter case, the toroidal field undergoes a reversal in the outer region of the vessel (hence the name).

For the tokamak configuration a much larger externally applied field B_ϕ is required and as a consequence this configuration was originally regarded as inflexible and the plasma produced difficult to diagnose, being hidden behind relatively massive toroidal coils. However, Thomson scattering measurements on the Russian tokamak machine T3 by a joint Culham/Kurchatov team (Peacock et al, 1969) showed that tokamaks produced conditions much closer to those required for ignition than did the reverse-field pinch system and research was concentrated on the tokamak system.

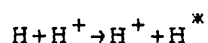
An idea of the progress towards breakeven (for the D-T reaction) may be obtained by seeing how the basic parameters have improved since 1969. Plasma currents have been increased from 150 KA to 5 MA, energy confinement times from 10 ms to 0.8 s and peak electron temperatures from 1 KeV to 3.5 KeV all on JET (Bickerton, 1985). The JET project located at UKAEA Culham Laboratory is the largest single project in the EURATOM fusion programme and when operating at its full design level should produce typical plasma currents of up to 4.8 MA, having a greater performance capability than any other machine of its type.

5.3 The Importance of Neutral Hydrogen in Fusion Plasmas.

The plasma is in close proximity to its containing vessel walls, and low density, cold neutral atoms will exist there. These atoms mostly originate from the dissociation of H_2^+ :



The neutrals produced will have an energy of a few eV and may then diffuse into the main body of the plasma or undergo charge exchange with protons to produce high energy neutrals:



These high energy neutrals (with a typical energy range in HBTX1A of 100eV $\rightarrow$ 2KeV (Bunting et al, 1985) will penetrate the plasma further. The neutrals will also be ionized by electron impact and the resulting neutral density profile will decrease with increasing distance from the walls.

These energetic neutrals may also hit the walls of the device resulting in the release of impurities (this

process is called sputtering). Hintz and Bogen (1984 and references therein) used fluorescence scattering on impurities to show that most of the iron impurities present in the TEXTOR tokamak were due to this process. Muller et al (1981) also reported that sputtering of the stainless steel walls from charge exchange neutrals appears to be the dominant production mechanism for impurity ions in the Doublet III tokamak.

Accumulation of the metallic impurities thus produced can give rise to enough radiation to seriously degrade the energy confinement (Duchs et al, 1977). The penetration of neutral beams used for auxiliary heating into the centre of the plasma may also be impaired by the presence of such impurities due to ionization by protons (>1 KeV) and charge exchange. Indeed, under certain conditions charge exchange between neutral beam atoms and high Z ions may dominate the overall power losses (Peacock, 1980). The presence of impurities also increase the cross section for ion impact ionization by approximately Z_{eff} (Girard et al, 1975) and this is important for the edge region if $Z_{\text{eff}} > 3$. Also, heating of the edge region may occur due to the deposition of the beam ions leading to increased ion thermal transport to the limiter, raising the typical energy of the charge exchange neutrals striking the wall. This would lead to greater sputtering and hence greater impurity density at the edge. To solve this problem, Duchs proposed the use of a cold plasma blanket to cool the edge region, thus depressing the energy of the charge exchange neutrals.

For increasing fractions of high Z impurities (such as tungsten, Mo, Fe, O) Jensen et al (1977) showed that the Lawson and ignition criteria increase. Note that a "clean" plasma with no impurities present is not ideal as regards ignition if the only heating mechanism available is ohmic heating. However, for additional heating (such as RF and neutral beam heating) impurity concentrations must be kept at a minimum. On JET for example, the use of

beryllium instead of the usual carbon limiters has been considered, to minimise Z_{eff} . The effects of impurities on fusion-type plasmas are discussed in more detail in Chapter 7.

5.4 Fluorescence Scattering as a Diagnostic.

Observation of the increase in signal (at the same wavelength) due to the laser pumping may be related to the initial level populations with knowledge of the relevant populating and depopulating rates between the levels. In the case of fluorescence scattering in the visible regions of the spectrum, i.e. the Balmer series, the enhancement is related to the ground state density by use of the 20 level CR code described by Gohil (1982) and referred to throughout this thesis. Fluorescence scattering as a diagnostic technique for the measurement of the absolute neutral density has several advantages over other methods such as Abel inversion due to the inherent spatial resolution available.

5.4.1 Resolution. The main advantages that fluorescence scattering as a diagnostic for fusion plasmas has over other techniques is the inherent spatial and temporal resolution available. The spatial resolution is limited simply by the irradiated volume. Under certain conditions it should be possible to obtain 1 cm resolution and detect a ground state density of 10^8 cm^{-3} for a signal to noise ratio of 2 (Gohil, (1982) and Gohil and Burgess (1983)). Similarly, the temporal resolution is fixed by the laser time duration (in this case FWHM=300 ns) and no complicated gating techniques are required so long as the detector is not saturated by the initial pinch emission. This may not be such an advantage on long lived plasmas although the time duration of the conventional dye laser used in this experiment may be increased to a few microseconds by altering the reflectivity of the output mirror.

5.4.2 Dependence on Radial Profiles. For temperatures greater than approximately 30 eV, the results from fluorescence scattering have minimal dependence on the exact radial temperature profile chosen. This is because for these high temperatures, the relevant rate coefficients between the levels do not change very much, the dominant contribution to the integral $\int \sigma(v) v f(v) dv$ being from regions where the cross section $\sigma(v)$ is relatively slowly varying. Given a ground state density of 10^{10} cm^{-3} , output from the CR code shows that for an increase in temperature from 50 eV to 100 eV (a factor of 2) for a typical electron density of 10^{13} cm^{-3} , the absolute enhancement Δn_3 of the upper pumped level will only increase from 7.8×10^6 to $8.7 \times 10^6 \text{ cm}^{-3}$ (a factor of 1.2). A further increase in temperature to 500 eV gives an absolute enhancement Δn_3 of $7.5 \times 10^6 \text{ cm}^{-3}$. Other line integrated techniques have a strong dependence on $T_e(r)$.

5.4.3 Other Methods. Other methods for obtaining the neutral density do exist but compromise on one or both of the above points. Abel inversion is the most widely used alternative to fluorescence scattering (as demonstrated on the TFR tokamak with the D- β line Equipe, TFR (1975) and also by Yamaguchi and Hirano (1981)). Abel inversion however requires very accurate knowledge of $n_e(r)$ and $T_e(r)$ since the numerical inversion techniques used are very sensitive to errors in the input data. In particular, problems arise from the contribution of the edge region due to the rapid change of plasma conditions there. Further, the line emission may not always be uniform and axial symmetry may not exist. Also, as tokamaks become larger, the geometric constraints will become more important limiting the number of profiles that may be obtained.

Barnett and Ray (1972) analysed the neutral density energy spectrum (usually used to measure the ion temperature) to obtain a neutral density profile.

Geometric considerations could adversely affect this method however. For example, for a plasma where $n_e > 10^{14} \text{ cm}^{-3}$ and minor radius $> 1 \text{ m}$, the plasma will be "thick" to neutrals (Wharton, 1975) and the neutrals from the centre would be unable to reach the spectrum analyser. Also, the transport of the neutrals from the centre to the edge requires modelling for the analysis of the energy spectrum.

5.5 Overview of Previous Work.

The first attempt to obtain a spatial measurement of neutral hydrogen using Balmer-alpha fluorescence scattering was by Skinner (1974) on the CLEO tokamak at the UKAEA Culham Laboratory, following an original suggestion by Burgess (1970). This experiment failed to produce a result primarily for lack of experimental time on the CLEO device. The first successful experiment with this technique however was performed on the FT1 tokamak device in the USSR (Razdobarin et al, 1979). Since then several other spatial measurements have been performed using the Balmer-alpha transition (see for example Reinhold et al 1982, Muraoka et al 1985).

5.5.1 Lyman-Alpha Fluorescence. All the experiments in hydrogen have been performed using the Balmer-alpha ($B-\alpha$, $n=2 \rightarrow 3$) transition in the visible region. It has been suggested by several authors (for example Breton and Papoular (1975), Koopman et al (1978), Steuer and Wrobal (1978), Bogen and Lie (1980)) that measurement of the ground state population via Lyman-alpha ($n=1 \rightarrow 2$) fluorescence scattering at $\lambda=1252 \text{ \AA}$ would be possible enabling a direct probing of the $n=1$ level and giving an increased enhancement due to the higher ratio of the initial level populations (for example $n_1^0/n_2^0=55$ for $n_e=5 \times 10^{14} \text{ cm}^{-3}$, $T_e=1 \text{ KeV}$). However, fluorescence scattering at this wavelength has many disadvantages (Gohil and Burgess, 1983). The main problem is the large powers required to saturate the $1 \rightarrow 2$ transition. Under

typical tokamak conditions the power required is approximately $800 \text{ kWcm}^{-2}\text{\AA}^{-1}$ (Gohil and Burgess, 1983). No Ly- α lasers have yet been developed capable of achieving this power density and any fluorescence scattering experiment carried out at this wavelength would have to be an unsaturated experiment. Therefore it would be necessary to know precisely not only the local laser power but also the homogeneity of the laser beam.

The long pulse duration produced by flashlamp-pumped lasers (up to 1 μs) used in Balmer-series fluorescence experiments allows for an improvement to the signal/noise ratio of the fluorescence signal simply by time-integrating the signal. Further, the high saturation powers required at Ly- α implies that much smaller volumes of plasma can be saturated at Ly- α than Balmer- α due to the limited laser powers available at Ly- α . Hence although Balmer-series fluorescence has a theoretically smaller signal/noise ratio than Lyman-series fluorescence (due to the smaller ratio of initial (pre-laser) level populations) the longer pulse lengths available and the easily attainable saturation intensities required make Balmer-series saturated fluorescence a much more feasible proposition than saturated Lyman-series fluorescence.

The experimental and interpretative difficulties of an unsaturated Ly- α experiment are much greater than a B- α saturated experiment especially since B- α lasers with sufficient saturation powers to are readily available. For a more detailed comparison between fluorescence scattering at Lyman and Balmer-alpha, see Gohil and Burgess (1983).

5.5.2 Fluorescence of Other Species. Fluorescence scattering has been applied to other species, for example neutral or low charge state impurities such as FeI, FeII, TiI and NiI (Muller et al, 1979, 1980, Cook et al, 1984) and NaI (Lipson, 1979). Lipson also measured the ion temperature with this technique by tuning the laser

across the Doppler width of the probed transition. Fluorescence scattering as a technique for a density measurement has yet to be exploited as regards highly ionized species. If this could be done (a study on one particular impurity is presented in Chapter 7) then it may be possible to use laser induced fluorescence of such ions as a localised diagnostic for the plasma interior.

5.6 Other Applications of Fluorescence Scattering.

As a diagnostic, laser induced fluorescence need not be restricted to neutral hydrogen density determination.

The application of this technique to the measurement of electron density and temperature was first discussed in the literature by Measures(1968). Dimock et al (1969) suggested that plasma drift motions could be measured with this technique, by observing the spatial variation in the fluorescence signal whilst pumping on two transitions. They measured the ion drifts in barium doped plasmas with densities as low as $2.5 \times 10^3 \text{ cm}^{-3}$ being measured. Burgess (1970) proposed a method to measure the magnetic field direction using this technique based on Hanle's third effect (Mitchell and Zemansky, 1934) and this is discussed in detail by Skinner (1974). Fluorescence scattering has also been used to measure the collision rates between excited states in helium. However, work of this kind relates to the low temperature regime which has been discussed in Chapter 1.

Impurity ion concentration may also be measured with this technique and direct fluorescence scattering on a given impurity line has been briefly discussed in Section 5.5.2. Gohil(1982) suggested an alternative method which involved pumping on and saturating the Balmer-gamma ($n=2 \rightarrow 5$) transition in hydrogen which would deplete the $n=1$ level (since this is strongly coupled to the higher levels under typical tokamak conditions). A decrease in the neutral density ground state would affect the charge

exchange rates (with impurity ions), reducing the abundance of the impurity lower charge states. Observation of the behaviour of a given ionisation state would give information about the density of the ion stage below. A similar type of experiment has been performed by Afrosimov (1978) but using an 8 KeV hydrogen beam rather than a laser to perturb the system, to estimate the density of C VII.

A spatial measurement of the electron density can also be obtained using the fluorescence scattering technique as described by Gohil (1982), and this is discussed in Chapter 6, section 6.4.3.

Chapter Six

FLUORESCENCE SCATTERING EXPERIMENTS ON A HIGH TEMPERATURE PLASMA.

In this chapter, a spatially resolved determination of the neutral deuterium density in HBTX1A by fluorescence scattering is described. The first set of measurements was made with a spatial resolution of 4 cm, the laser being tuned to the Balmer- α (6561 Å) transition. The second set was performed with a resolution of both 4 cm and 2 cm, with and without the vertical field, the laser being tuned to the Balmer- β (4860 Å) transition. Improved signal to noise was obtained by moving to this wavelength. A slightly asymmetric neutral profile was observed, the plasma not being central in the vessel. Addition of a vertical field produced a reduction in the neutral density, this reduction being more pronounced towards the plasma centre. For comparison, neutral density profiles were generated using the AURORA code, and the effects of different radial density and temperature profiles were investigated.

6.1 Experimental Requirements.

The experimental technique of laser induced fluorescence scattering when applied to high temperature fusion type plasmas is much the same as when applied to low temperature z-pinch plasmas, as described in Chapter 1.

Uncertainties regarding the electron collision rates in plasmas where T_e is close to the threshold of the Balmer- α transition (discussed in Chapter 2) do not affect the analysis of data in the high temperature case since all the available theories tend towards the Born approximation.

The main points that need to be considered when employing fluorescence scattering techniques under fusion plasma conditions are:

6.1.1 Ground State Coupling. In HBTX1A the electron temperature T_e (100 eV at the plasma centre) is much larger than both the transition energy between the ground and the first excited state ($E_{12}=10.2$ eV) and the ionization energy of deuterium ($I_p=13.6$ eV). The ground state is therefore coupled to the excited states and the continuum via electron collisions. If a particular transition in the visible region (e.g. a member of the Balmer series) is pumped and saturated, the observed fluorescence enhancement at this wavelength is therefore related to the ground state density because of this coupling. This is in sharp contrast to the work on the low temperature (<1 eV) plasmas previously discussed, where the ground state is effectively decoupled from the excited states and so these plasmas are not suitable for such a diagnostic. The absolute enhancement is related to the ground state population by use of the 20 level CR code as described by Gohil (1982) and referred to throughout this thesis, which incorporates the relevant atomic processes.

6.1.2 Power Required for Saturation. The power required to achieve saturation for the D- α (6561 Å) and D- β (4860 Å) transitions for temperatures of 100, 500 and 1000 eV as a function of electron density are shown in Figure 6.1 (taken from Gohil (1982), calculated from Equation (1-2)). Note the linear dependence on the electron density of the D- β saturation power. This is because the collision rates out of the $n=4$ level to the higher levels are much larger than the radiative rate A_{42} . For the D- α transition the power is virtually independent of the electron density for $n_e < 10^{13} \text{ cm}^{-3}$, since the radiative rate A_{32} is much larger than the collision rates associated with the $n=3$ level for this set of plasma conditions. Note also that larger powers are required to saturate the D- β transition by approximately an order of magnitude. This is because of the shorter wavelength of the transition, the lower radiative rate A_{42} and because the collision rates out of level 4 are larger than those out of $n=3$ under these plasma conditions. For HBTX1A densities (10^{13} cm^{-3}) the saturation powers required for the D- α transition are within the specifications of conventional flashlamp pumped dye lasers. Previous work involving a dye laser operating at the D- β wavelength did not produce sufficient power for saturation under these conditions (Ward, 1983).

6.1.3 Enhancement Expected. Figure 6.2 shows the temporal behaviour of each of the upper level populations $n=3$ and $n=4$ resulting from saturation of the D- α and D- β transitions. This was generated using the CR code for typical HBTX1A conditions of $n_e = 10^{13} \text{ cm}^{-3}$ and $T_e = 100 \text{ eV}$. The flat-topped, square shaped laser pulse used had a FWHM of 300 ns, a risetime of 10^{-14} s and a spectral bandwidth equal to 1.5 times the Doppler width. Similar to the 1 eV (z-pinch) plasmas, the enhancement is independent of the laser intensity once saturation has been achieved. However, in contrast to the low temperature case, the characteristic transient spike is no longer seen. The fluorescence signal rises to a

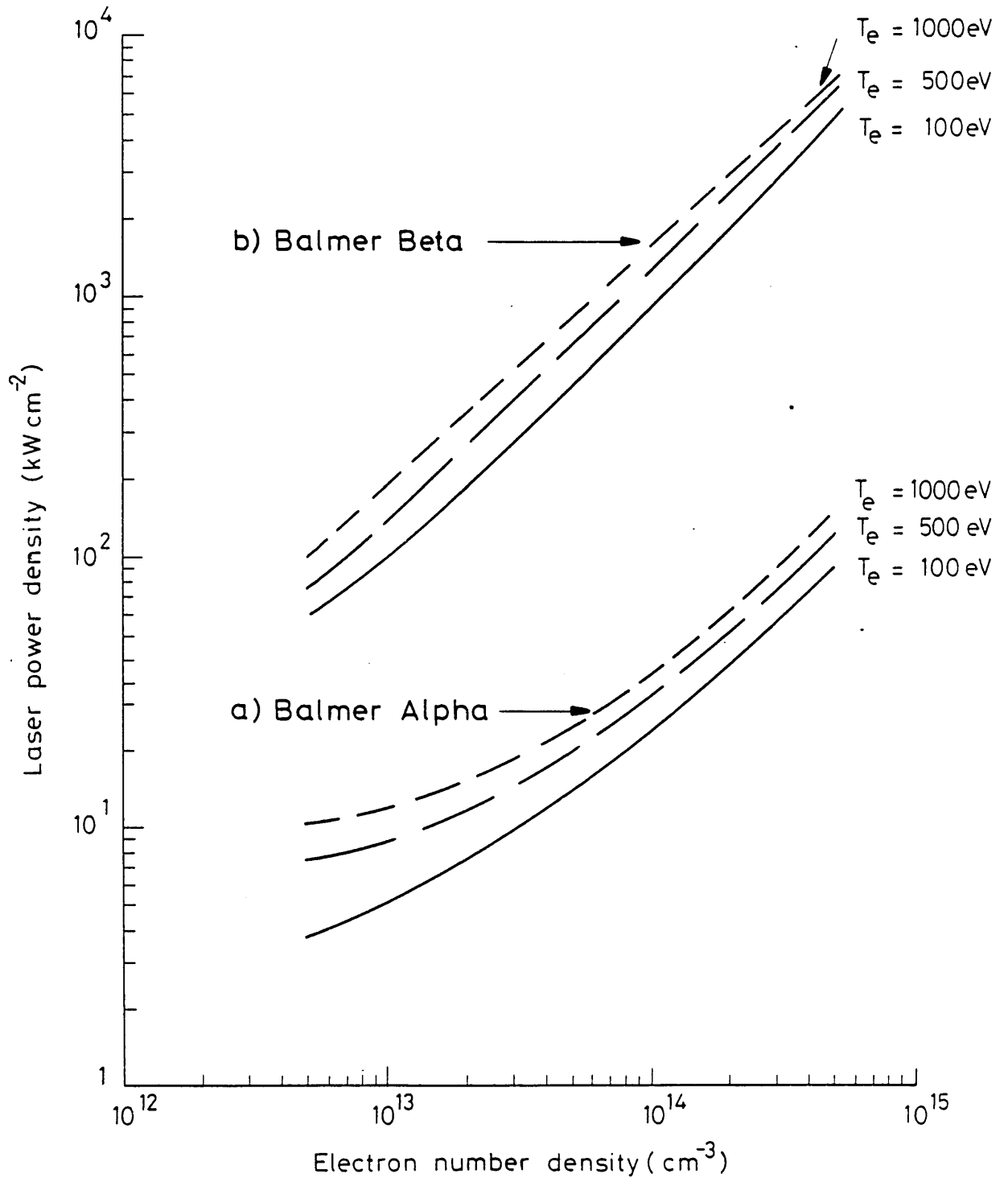
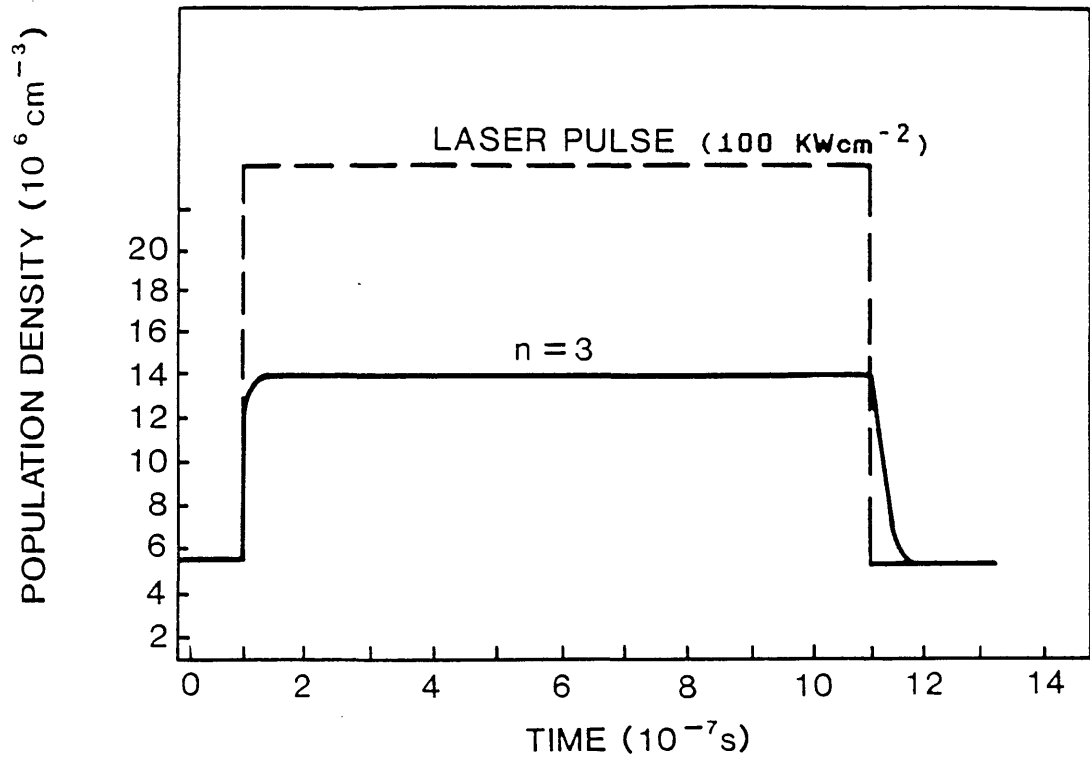


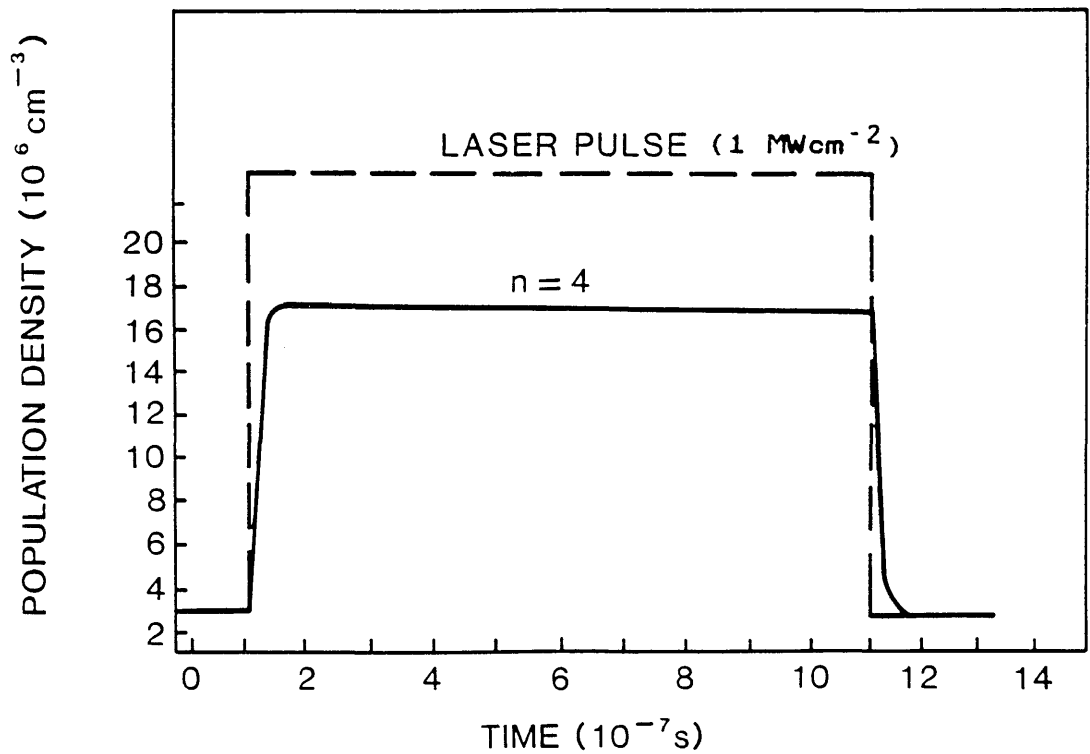
Figure 6.1 The Laser Power Density Required to Saturate the Balmer-Alpha and Beta Transitions.

This shows the theoretical saturation powers required for the H- α and H- β transitions (into 1.5 times the Doppler width) in high temperature plasmas for a saturation factor $f_g=0.75$.

Figure 6.2 The Temporal Behaviour of the $n=3$ and 4 Level Populations Whilst Saturating the Balmer-Alpha and Beta Transitions Respectively.



1) The $n=3$ Level Behaviour (H-Alpha Saturation).



11) The $n=4$ Level Behaviour (H-Beta Saturation).

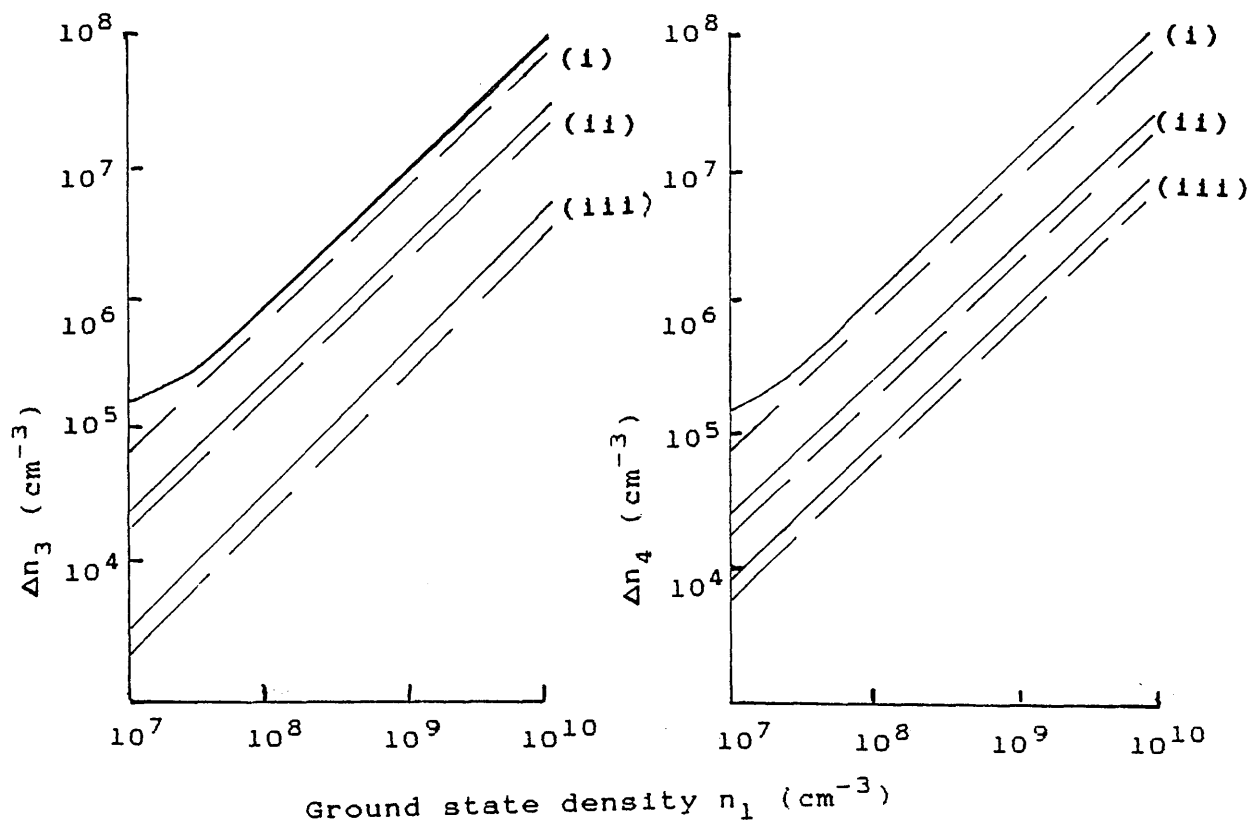
For both cases, the ground state density is taken to be 10^{10} cm^{-3} and $T_e=100 \text{ eV}$, $n_e=10^{13} \text{ cm}^{-3}$.

constant value as the laser pulse turns on. Note that once the laser pulse is terminated, the population relaxes to its initial value with a time given by the dominant rate out of the upper level. This fact may be used as a diagnostic for the local electron density which will be discussed in Section 6.4.3.

Figure 6.3 shows the actual change in the upper level population Δn_3 , Δn_4 whilst saturating the D- α and D- β transitions, as a function of the ground state density n_1 . Results from the CR code are shown for a range of electron densities and temperatures. The linear dependence of Δn_u on n_1 is clearly shown (this is due to the strong coupling between the $n=1$ level and the excited states as previously explained). Knowledge of the absolute enhancement of the upper level Δn_u will then give the $n=1$ level population directly. This linear dependence breaks down for a low n_1 population ($n_1 < 10^7 \text{ cm}^{-3}$), high electron density ($n_e > 5 \cdot 10^{14} \text{ cm}^{-3}$) and the low temperature ($T_e < 100 \text{ eV}$). Under these conditions, the collision rates associated with the $n=3$ and 4 levels up to the higher levels are much larger than the collision rates down to the ground level. Note also that Δn_3 and Δn_4 are relatively insensitive to changes in temperature under these conditions as briefly discussed in Section 5.4.2.

6.1.4 The Expected Signal to Noise Ratios. The signal to noise ratios expected for both D- α and D- β fluorescence scattering are discussed in some detail by Gohil (1982). For typical HBTX1A conditions of $n_e = 10^{13} \text{ cm}^{-3}$, $T_e = 100 \text{ eV}$ and a neutral density $n_1 = 10^{10} \text{ cm}^{-3}$ a signal to noise ratio for D- α of 5.5 and for D- β of 7.6 would be expected. These figures were obtained assuming a solid angle of $2 \cdot 10^{-3} \text{ Sr}$, an irradiated volume of 1 cm^3 (background volume 100 cm^3) and a detection efficiency of 5% at the D- α wavelength and 10% at the D- β wavelength.

a) Balmer-Alpha Saturation. b) Balmer-Beta Saturation.



KEY:

- - -	$T_e = 1000$ eV	(1)	$n_e = 5 \cdot 10^{14}$ cm ⁻³
—	$T_e = 100$ eV	(11)	$n_e = 5 \cdot 10^{13}$ cm ⁻³
		(111)	$n_e = 5 \cdot 10^{12}$ cm ⁻³

Figure 6.3 The Absolute Change in the Upper Level Population due to Saturation as a Function of the Ground State Density.

The data shown is from Gohil (1982) and shows Δn_U , the absolute change in the upper level population for saturating the H-Alpha and H-Beta transitions.

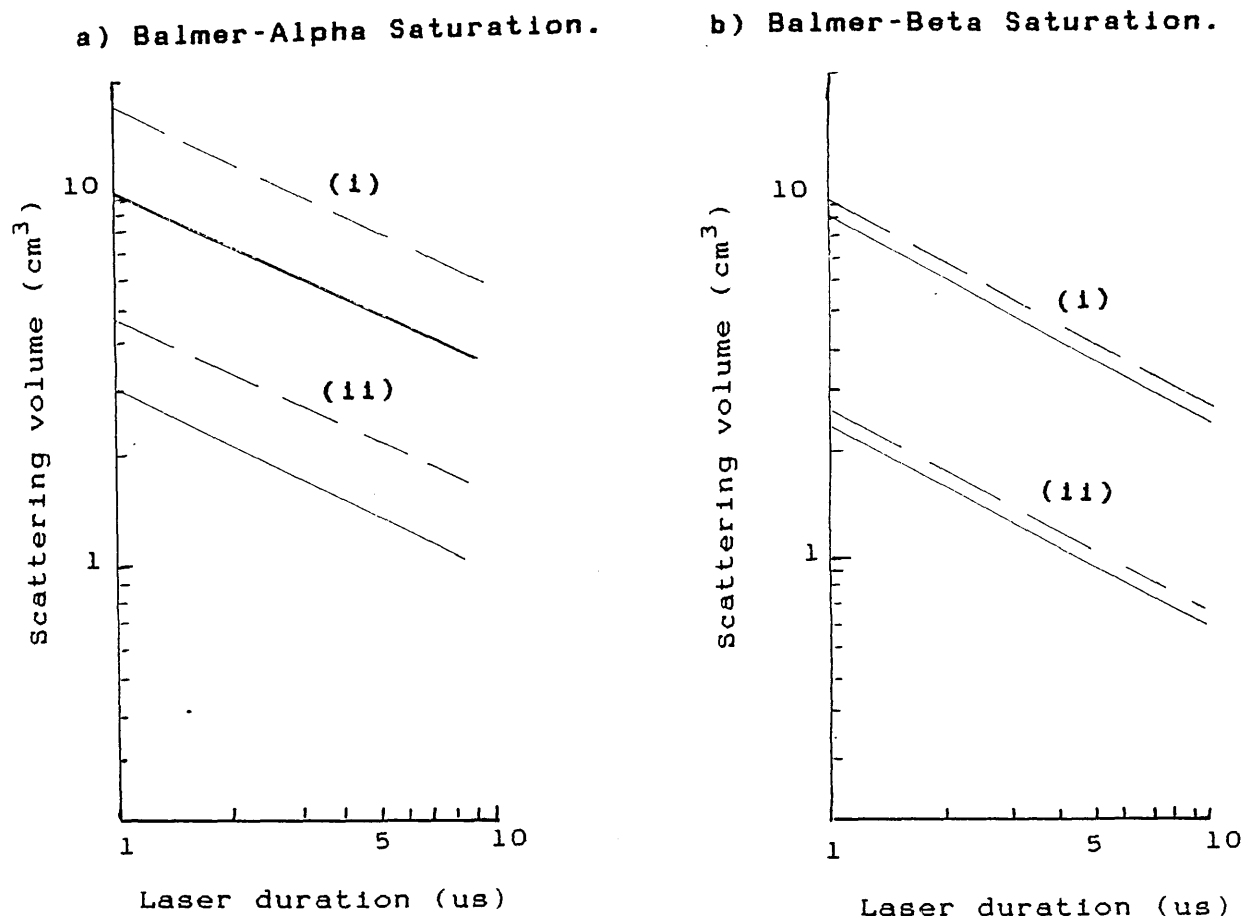
6.1.5 Sensitivity. The "sensitivity" of the fluorescence technique can be defined as being the minimum ground state density n_1 for which the signal to noise of the resulting fluorescence signal is equal to 2 (Gohil, 1982). Figure 6.4 shows the minimum n_1 density that could be detected by this criterion as a function of the scattering volume and laser time duration for a given electron density with the same assumptions as in Section 6.1.4. Given the right combination of plasma conditions and detection system, neutral densities as low as 10^7 cm^{-3} should be detectable, depending on the spatial resolution required. There is a slight advantage as regards increased sensitivity for the D- β transition because of the greater change Δn_4 in the upper level population during saturation at the lower electron densities ($n_e < 5 \cdot 10^{13} \text{ cm}^{-3}$). But for $n_e > 5 \cdot 10^{13} \text{ cm}^{-3}$, D- α fluorescence scattering offers the better sensitivity. Note also that the sensitivity decreases as the temperature increases, due to the lower $n=2$ density prior to the laser irradiation.

6.1.6 The Required Spectral Bandwidth. The laser spectral bandwidth should ideally be larger than the linewidth of the transition studied so that all the available neutral atoms are excited to give the greatest fluorescence signal. This linewidth contains the following contributions:

1) Doppler Width. The Doppler width of a spectral line (FWHM) is given by the following expression:

$$\Delta\lambda = \frac{2\lambda}{c} \left(\frac{2(\ln 2)KT_e}{m_H} \right)^{1/2} \quad (6-1)$$

where T_e is the electron temperature (=neutral hydrogen temperature) and M_H is the mass of the hydrogen atom and all other symbols have their usual meanings. For typical HBTX1A operating temperatures of 50-100 eV, this corresponds to a range of 2.5-3.5 Å for the D- α transition and 1.8-2.2 Å for the D- β transition.



KEY:

(1) $n_1 = 10^7 \text{ cm}^{-3}$

(11) $n_1 = 10^8 \text{ cm}^{-3}$

Figure 6.4 The Minimum Ground State Density Resolvable as a Function of the Scattering Volume and Laser Duration.

This shows the minimum n_1 that can be detected (for an assumed signal to noise ratio of 2, see text) and is taken from the calculations of Gohil (1982) for $T_e = 100 \text{ eV}$ (—), $T_e = 1000 \text{ eV}$ (---) and $n_e = 10^{13} \text{ cm}^{-3}$.

ii) Stark Width. The Stark widths were taken from the tables of Vidal et al (1973). The maximum temperature at which the widths were tabulated was $4 \cdot 10^4$ K (4.65 eV). The data presented in Table 6.1 corresponds to this temperature and an electron density of $3.16 \cdot 10^{13} \text{ cm}^{-3}$.

iii) Zeeman Splitting. From Mitchell and Zemansky (1971) the Zeeman splitting $\Delta\lambda$ for two levels 1 and 2 is described by:

$$\Delta\lambda = 0.466\lambda^2 B (g_{J1}M_1 - g_{J2}M_2) \cdot 10^{-9} \quad (6-2)$$

where the wavelength λ is in units of Å, B is the magnetic field (=1 T for HBTX1A), g_J are the Lande g-factors and M_1 and M_2 are the magnetic quantum numbers for levels 1 and 2. The Lande g factors for LS coupling and weak magnetic fields are given by:

$$g_J = \frac{1 + j(j+1) + s(s+1) - l(l+1)}{2j(j+1)} \quad (6-3)$$

For a first approximation it is assumed that the factor $(g_{J1}M_1 - g_{J2}M_2) = 1$. It should also be noted that the different Zeeman components have differing absorption coefficients so the laser spectral bandwidth must be greater than the Zeeman splitting of the line to avoid this complication.

Table 6.1 shows the Doppler and Stark widths (FWHM) and Zeeman splitting for typical fusion-type plasma conditions for the first three Balmer transitions in deuterium. From the tabulated values it can be seen that the Doppler width has the much the largest contribution to the total linewidth in each case, and it is this value that the laser linewidth must exceed if possible.

TABLE 6.1: Comparison of the Doppler Widths, Stark Widths and Zeeman Splitting for the First Three Balmer Transitions in Deuterium.

	Doppler Width		Stark Width	Zeeman Splitting
	100 eV	1 KeV	($n_e = 3 \cdot 10^{13}$)	(B=1 Tesla)
D- α	3.5	11.3	0.06	0.20
D- β	2.6	8.3	0.20	0.11
D-g	2.3	7.4	0.02	0.09

This table shows several line broadening mechanisms relevant to the HBTX1A plasma, calculated for the first three members of the Balmer series of neutral deuterium, calculated for plasma conditions typical of the HBTX1A device (except for Doppler broadening at $T_e = 1\text{KeV}$ which is included only for comparison). It can be seen that Doppler broadening is the most important mechanism. The widths are all quoted in Angstroms.

6.1.7 Opacity. For Doppler broadened spectral lines, the optical depth t at line centre is given by Griem and Lovberg (1970) to be:

$$t = (e^2/mc^2) N \lambda_0 f d (\pi m c^2 / 2KT)^{1/2} \quad (6-4)$$

where m is the atomic mass, N is the absorber density, λ_0 is the line centre wavelength ($\text{\AA}$), f the oscillator strength of the transition and d is the path length. All other symbols have their usual meanings. The oscillator strength of a transition between two levels i and j is given (Thorne, 1974) by:

$$f_{ij} = 1.499 \cdot 10^{-16} (g_j/g_i) \lambda_0^2 A_{ji} \quad (6-5)$$

where A_{ji} is the spontaneous decay rate (s^{-1}) of the transition and λ_0 is the peak wavelength in $\text{\AA}$.

For an optically thin transition $t < 1$. At 100 eV, the Nd product which gives unit optical depth is $2.9 \cdot 10^{17} \text{ cm}^{-2}$ for the Lyman- α transition and $3.9 \cdot 10^{16} \text{ cm}^{-2}$ for the Balmer- α transition. For a maximum $n=1$ density of 10^{12} cm^{-3} in HBTX1A (diameter 50 cm) the Nd product for the Lyman- α transition is $5.0 \cdot 10^{13} \text{ cm}^{-2}$ and for the Balmer- α transition is $5.0 \cdot 10^{10}$ (where the $n=2$ density is typically a factor of 10^{-3} smaller than the ground state density). Therefore all hydrogen lines in the UV and visible spectral regions in HBTX1A are optically thin, and the escape factors associated with these transitions were all set to unity in the CR code.

6.2 The Laser.

A xenon pumped coaxial dye laser, as described in Chapter 4 was used to probe the D- α and D- β transitions. The laser consisted of an oscillator amplifier arrangement. This is shown in Plate 6.1 along with the dye circulation unit. A degree of tunability was provided by a single Fabry-Perot etalon placed in the oscillator cavity as

shown. A different etalon was used for each transition. Both were manufactured to specification by Imperial College Optical Systems, having a 2.5 cm aperture, 5 μ m spacing and 70% reflectivity over the visible region, being flat to $\lambda/100$.

It was decided to improve upon the design of the laser heads. One of the problems associated with this type of laser was that the flash tubes needed to be changed frequently as damage was caused by the high voltage discharge. The outer tube was secured in place against the insulator by glueing it with a flexible, insulating adhesive. The adhesive used was silicon rubber. However, this took 24 hours to cure, which led to delays, particularly if the seal was not made carefully enough since the tubes would simply break again when subjected to voltages greater than 18 KV.

Several alternative fast setting adhesives were experimented with in an attempt to cut down on the time needed to replace the tubes. This was of particular importance for work at the D- β wavelength since consistently high powers were required to achieve saturation, and so the flash tubes needed to be cleaned and/or changed on a regular basis. Great care needed to be exercised in the dismantling and reassembling of the laser heads since any over tightening of the head would lead to broken tubes, the bores and tube diameters being very critical.

This aspect of the laser head was redesigned. A PTFE sleeve was made to hold the outer flash tube in place preventing high voltage from tracking across by increasing the distance over which the spark would have to travel. The modified laser head is shown in Figure 6.5. This simple modification worked very well, and made changing the flashlamps a much easier and less risky task.

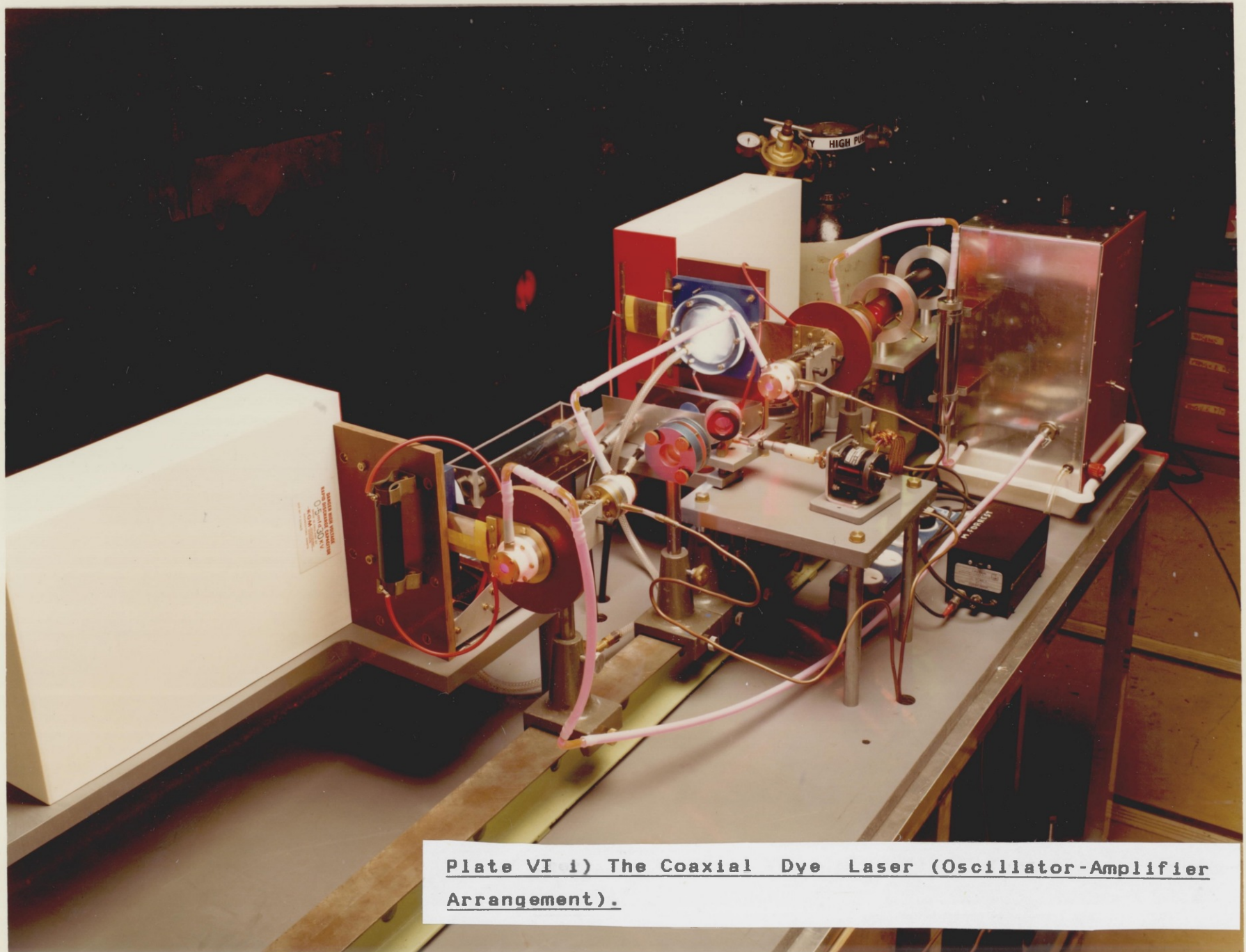
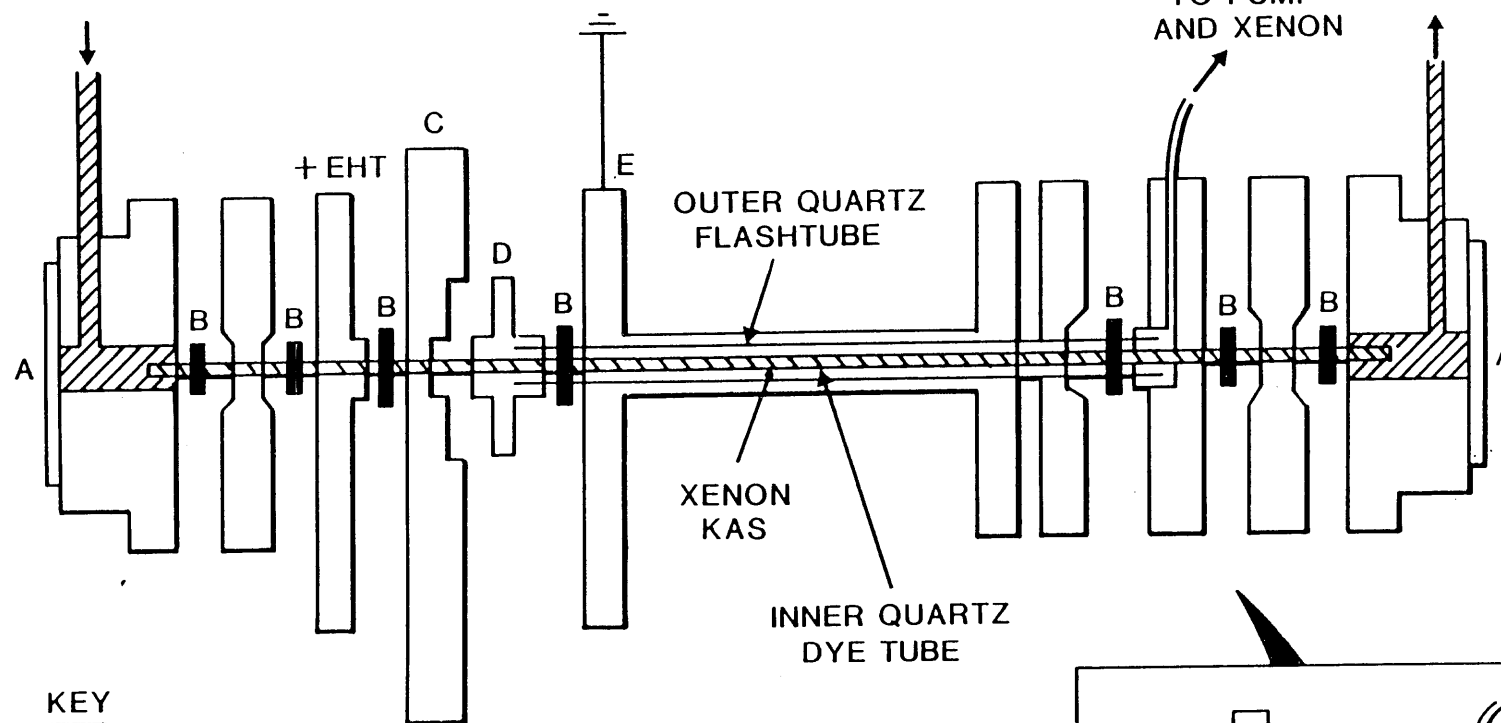


Plate VI 1) The Coaxial Dye Laser (Oscillator-Amplifier Arrangement).



A = END WINDOW

B = 'O' RING

C = PERSPEX INSULATOR

D = PTFE PIECE ADDED TO
PREVENT TRACKING

E = CENTRAL BRASS SECTION

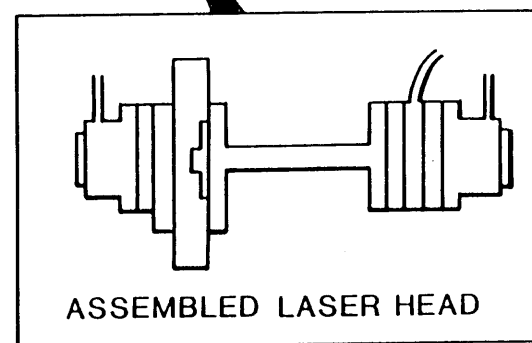


Figure 6.5 The Modified Dye Laser Head Showing Details of Assembly.

6.2.1 Lasing Action at the Balmer-Alpha Wavelength. For operation at the D- α wavelength ($\lambda=6561 \text{ \AA}$), a dye solution of 13 mg/l of cresyl violet and 26 mg/l of rhodamine 6G from Oriel Scientific Ltd was selected, the solvent being methanol. The back mirror was dielectric coated for 100% reflectivity and the output mirror for 50% reflectivity at this wavelength.

This arrangement typically gave 200 mJ output at a working voltage of 19 KV and a xenon pressure between the flash tubes of 80 mbar into a pulse of time duration 300 ns (FWHM) and measured bandwidth 5 $\text{\AA}$. This was an efficiency of 0.22% and a power of 493 KWcm^{-2} into the Doppler width (3.7 $\text{\AA}$ at 100 eV see Table 6.1).

6.2.2 Lasing Action at the Balmer-Beta Wavelength. For lasing at this wavelength ($\lambda=4860 \text{ \AA}$), the etalon, mirrors, dye solution and filter had to be changed, the whole system being thoroughly flushed through with clean methanol to remove any trace of the previous dye.

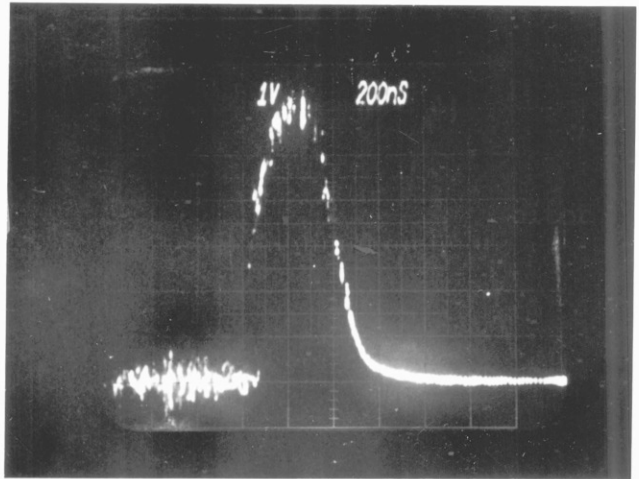
The dye solution used was Coumarin 102 from Oriel Scientific Ltd, methanol being the solvent. Previous work with the same dye at this wavelength (Ward, 1983) had not been successful, partly because of the original design of the laser head itself, which led to continual breaking of the flash tubes. Experiment showed that a dye concentration of 35 mg/l and a xenon pressure of 100 mbar (as opposed to the usual 80 mbar for D- α operation) gave the optimum power output whilst still retaining a good homogeneous laser pulse. This combination gave an energy output of typically 450 mJ at 19 KV (an efficiency of 0.50%). With a pulse duration of 610 ns (FWHM) this gave an output power of 734 KWcm^{-2} into the bandwidth (5 $\text{\AA}$) and 384 KWcm^{-2} into the D- β Doppler width (2.6 $\text{\AA}$ at 100 eV see Table 6.1).

It was noticed that the FWHM time duration of the laser pulse at the D- β wavelength was approximately 290 ns

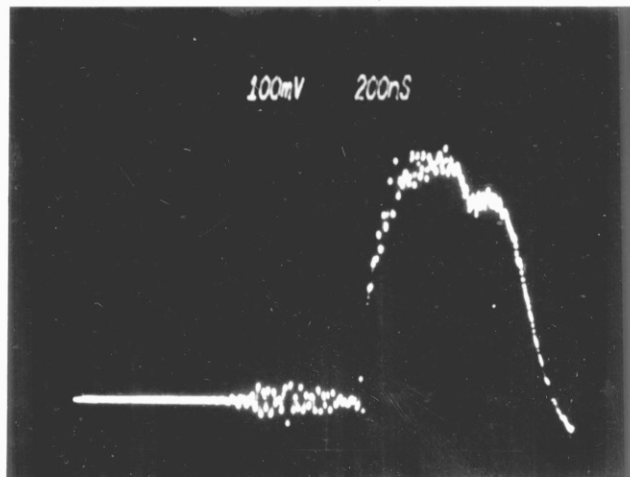
PLATE VI ii)

Dye Laser Photodiode Traces

Balmer-alpha laser monitor
FWHM = 360 ns



Balmer-beta laser monitor
FWHM = 610 ns



longer than the D- α laser pulse. The reason for this was unclear since the D- β pulse duration did not change with xenon pressure or wavelength within the tuning range of the Fabry-Perot etalon (approximately 4760-4960 Å). Also, it was noticed that if the oscillator was fired by itself (i.e. without the amplifier) the pulse displayed the same time history as the oscillator/ amplifier arrangement and so this effect was not due to the amplifier. Plate VI 11) shows a typical laser pulse for both the D- α and the D- β wavelength displayed on a Tektronix 7904 oscilloscope and detected by an EMI 9658A photomultiplier. The FWHM width of the D- α laser pulse was 320 ns and of the D- β laser, 610 ns.

The mirrors for the D- β transition were coated for 100% and 50% reflectivity at this wavelength and were flat to $\lambda/10$.

In order to maintain the high power output from the D- β laser system adjustments had to be continually made. Fresh dye was required (lasting approximately 300 shots) and the xenon changed frequently (every 20 shots or so). The quartz flash tubes were also changed or cleaned weekly.

6.3 Experimental Arrangement.

Figure 6.6 shows a schematic illustration of the experimental set up for the scattering arrangement. Plate VI 111) shows a photograph of the laser firing into HBTX1A. Note that the detection system is situated on top of the device, looking down into the plasma.

6.3.1 Input Optics. The laser was focussed through a baffle system into the plasma vessel, and was dumped into OBA blue glass to reduce the stray light. The input baffles consisted of a three baffle system of circular blue glass stops, the input baffle being 30 mm in diameter and the output baffle 20 mm, the entire system

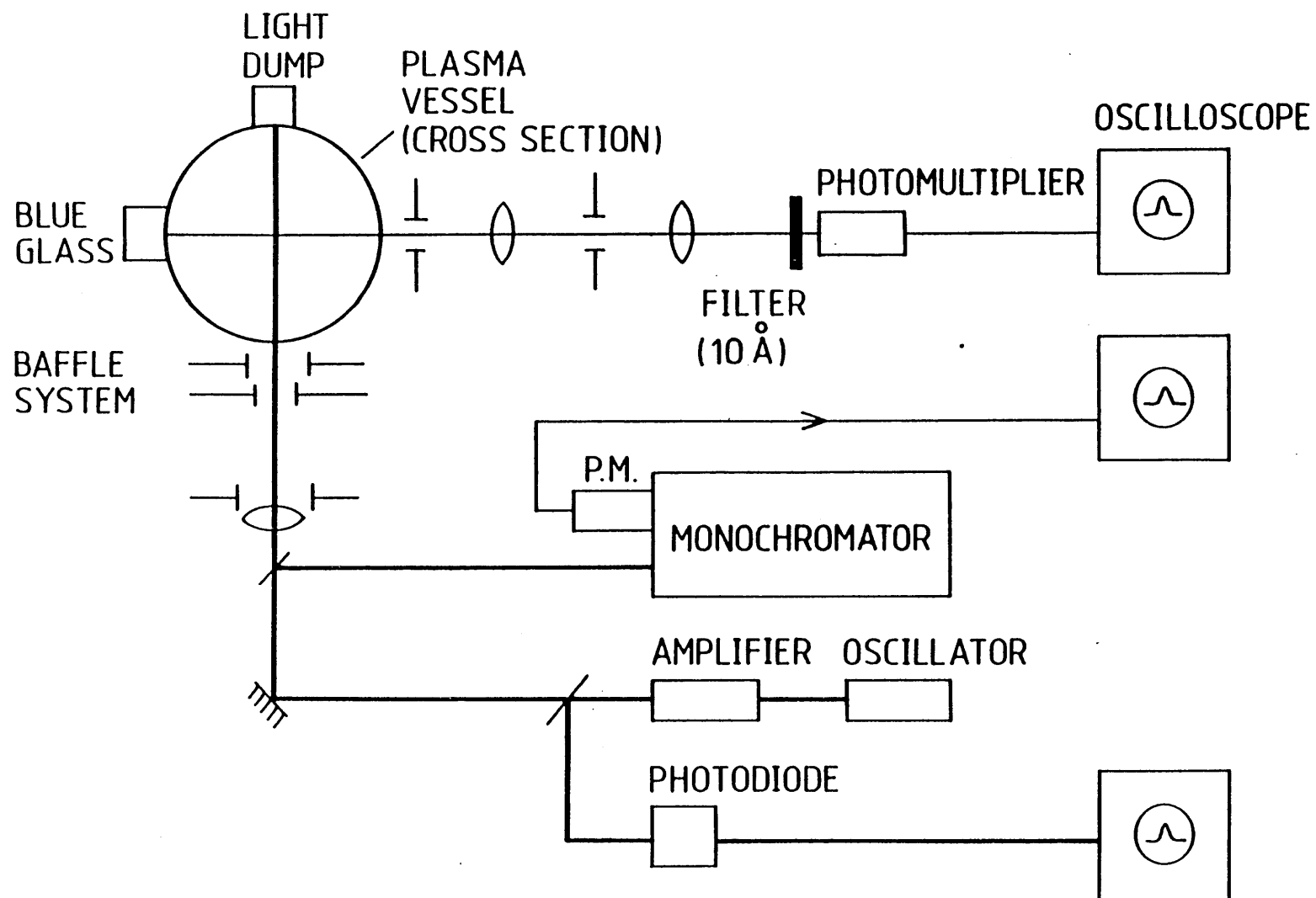


Figure 6.6 The Experimental Arrangement for Fluorescence Scattering on HBTX1A.

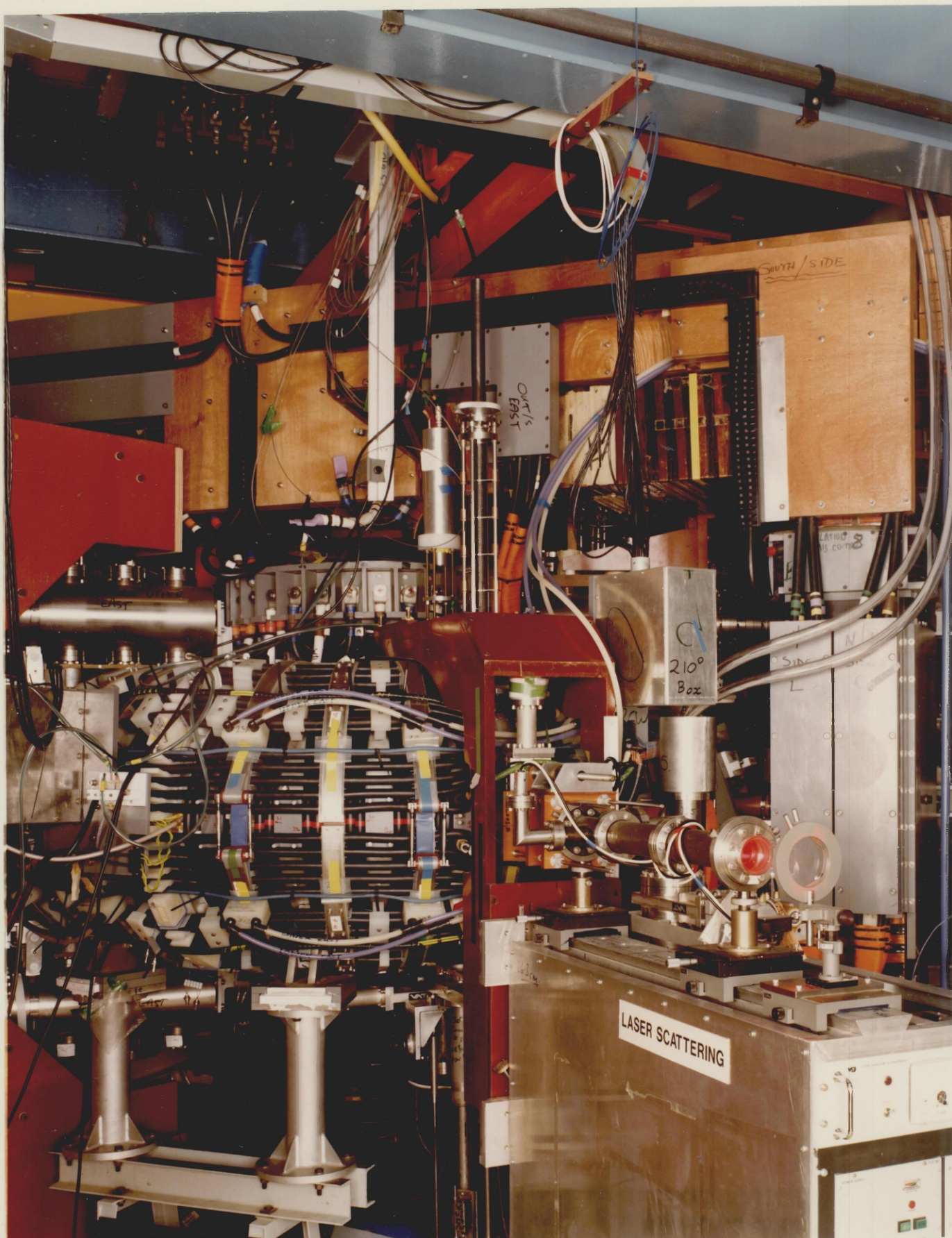


Plate VI 111) HBTX1A With the Balmer-Alpha Laser Firing
onto the Entrance Baffle System.

being 1 m long. This arrangement was used partly to reduce the stray light level but also to make the fluorescence scattering experiment compatible with the Thomson scattering system already in existence which used the same input window. For the two systems to be used in tandem, two 5 cm diameter mirrors were specially coated. The first was made to have 100% reflectivity at 45° incidence for the D- α wavelength and the other to have 100% reflectivity at the D- β wavelength, both allowing 100% throughput for the ruby light. To change to Thomson scattering mode, it was only necessary to remove the convex lens that focussed the dye laser into the baffle system, ensuring that alignment was not lost for either diagnostic.

6.3.2 Output Optics. The fluorescence signal was viewed at 90° to the input beam. A schematic illustration of the output optics is shown in Figure 6.7 and described below.

The optical system consisted of firstly a 20° wedge mounted in a circular 3 cm diameter stop that was able to be rotated in order to view the relevant positions in the plasma (a total of three different positions being possible for a given output window). The light then passed through a telescope arrangement of 15:1 demagnification, a rectangular slit which defined the viewed volume and a plano-convex lens of focal length 15 cm which produced parallel light. In the parallel light was situated an interference filter and another plano-convex lens of focal length 15 cm to focus the light onto the photocathode of an EMI 9658A photomultiplier. The interference filter was used as a monochromator to distinguish the line emission from the background. The D- α filter had a spectral bandwidth of 10 Å centred at 6561 Å and the D- β filter had a 5 Å bandwidth centred at 4860 Å. The rectangular aperture had dimensions 5.2x2.2 mm and the system was such as to image the slit into the plasma at the laser position, effectively magnifying the slit 15 times and defining the

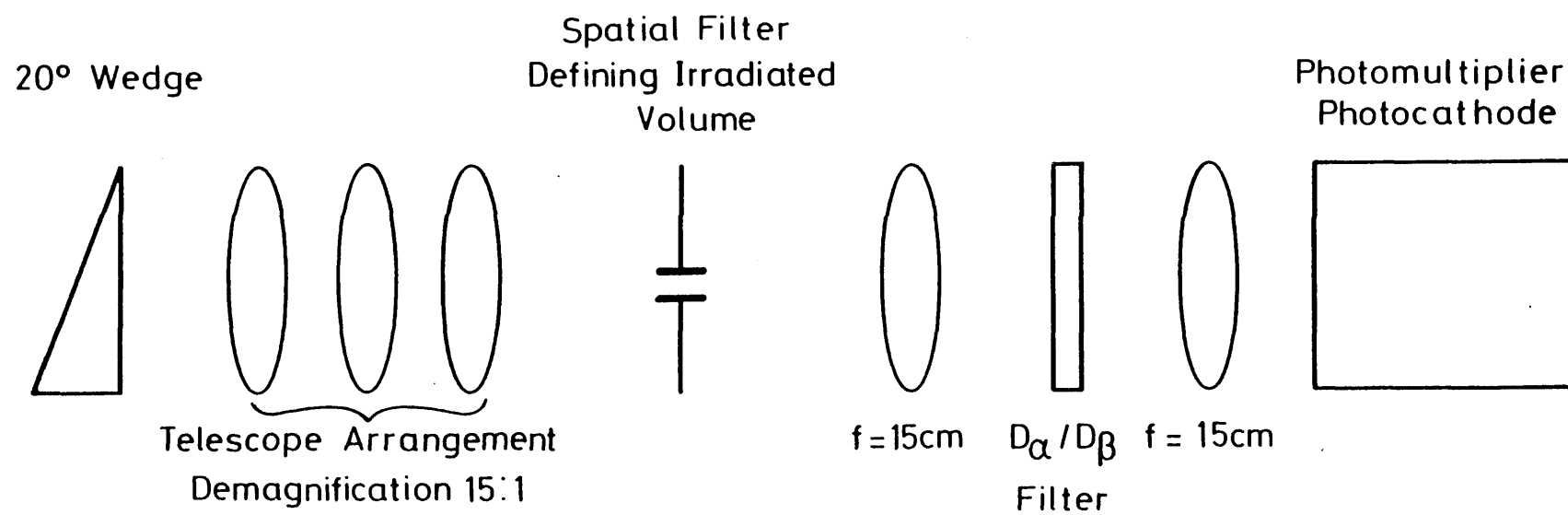


Figure 6.7 The Detection Optics.

volume under investigation. For the position 9 cm from the inner wall limiter (see Figure 6.11 for the geometric details) the irradiated volume was 58.58 cm^3 . Given operating conditions of $T_e=100 \text{ eV}$, $n_e=3 \cdot 10^{13} \text{ cm}^{-3}$ this gave an approximate limit of 10^8 cm^{-3} as the minimum neutral density detectable (see Section 6.1.5). The output was finally displayed on a Tektronix 7904 oscilloscope (7A24 plug-in) and photographed on 612 polaroid film.

6.3.3 Laser Monitor. The laser pulse was monitored via a fast ITL photodiode (risetime $<1 \text{ ns}$) that was placed to view the reflections of the laser from the Fabry-Perot tuning etalon. A beam splitter was placed in the path of the laser beam and the laser was also monitored via a Hilger 1000 monochromator and an EMI 9816 photomultiplier. The output signals were displayed on 7904 oscilloscopes.

6.3.4 Timing. For both the D- α and the D- β fluorescence scattering experiments the laser was triggered 1.5 ms after the peak of the observed emission. The plasma was reproducible at this time, the current in the coils being flat topped. The shot to shot reproducibility of the D- α emission line was better than 15%, and that of the D- β better than 10% at this delay time. Plate VI iv) shows a typical shot of the D- α and D- β emission, as seen through the output optics used in this experiment, for the duration of one HBTX1A discharge. The laser jitter for each wavelength was less than 10 ns, this being negligible compared to both the plasma lifetime (3 ms) and the time scale over which plasma conditions changed.

6.3.5 Check for Saturation. It was necessary to perform checks to confirm that the signal observed was really due to fluorescence scattering and not stray light. For each transition, the laser was tuned through the line and the variation in the signal observed. This varied as would be expected, being effectively zero outside the line width

and reaching a maximum at line centre. Thus the scattering was a resonance phenomena, and was not due to scattering from a dust particle. The level of stray light was also checked after each fluorescence shot. This was defined as the signal obtained when firing only the laser into the vessel (no plasma), and was negligible compared to the fluorescence signal. For each fresh fill of xenon, saturation of the probed transition was checked at the beginning and end of each fill by placing a 0.3 gelatine neutral density filter in the laser path (reducing the laser power by a factor of $10^{0.3}=2$). If the fluorescence signal observed retained the same magnitude and time history, then the laser was considered to be saturating the transition. On an individual shot-to-shot basis, a good indication that saturation had been achieved was to observe the risetime of the fluorescence signal. If the signal rose to its peak value faster than the laser pulse then this was also an indication of saturation.

6.3.6 Calibration. For absolute calibration of the system, an NPL standard irradiance lamp (quartz iodine glass cresta source) was used. The output of the lamp was reflected onto a large (100x100 mm) opal flat, the light being reflected into 2π . The spectral irradiance quoted for the D- α wavelength was $4.513R \text{ mWm}^{-2}\text{nm}^{-1}$ and for the D- β wavelength $13.7R \text{ mWm}^{-2}\text{nm}^{-1}$ where R is the reflectivity of the opal flat at that wavelength.

The absolute enhancement Δn_U of the upper level was calculated from the equation:

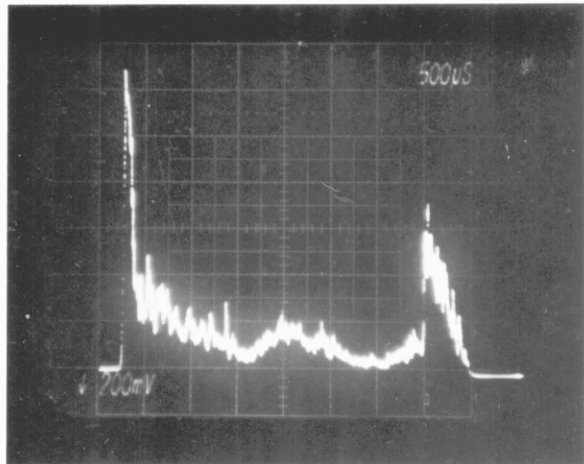
$$\Delta n_U = (4\pi F_R / V_F \Delta\Omega A_{UL}) \quad (6-6)$$

where F_R is the rate of fluorescence photons detected as calculated from the absolute calibration, V_F is the irradiated volume, $\Delta\Omega$ is the solid angle (10^{-2} Sr) and A_{UL} the spontaneous radiative rate s^{-1} of the transition considered ($A_{UL}=4.4 \cdot 10^7 \text{ s}^{-1}$ for D- α , $8.4 \cdot 10^6 \text{ s}^{-1}$ for D- β).

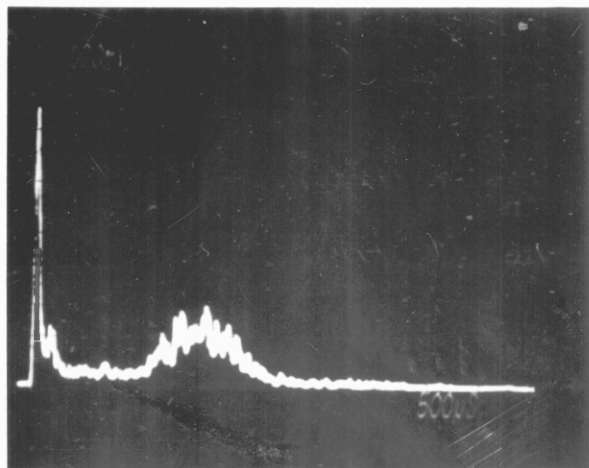
PLATE VI iv)

Balmer-alpha and Balmer-beta Emission in HBTXL1A

Balmer-alpha emission
500 mV/div 500 μ s/div



Balmer-beta emission
100 mV/div 500 μ s/div



The calibration factor for the system was calculated to be 5.82×10^{12} photons $s^{-1}mV^{-1}$ into 50Ω termination for the D- α transition (for a photomultiplier operating voltage of 800 V) and 3.86×10^{12} and 5.4×10^{12} photons $s^{-1}mV^{-1}$ into 50Ω for the D- β transition (for operating voltages of 800 and 850 V respectively and a neutral density filter of 0.6 in front of the photomultiplier). In the D- β fluorescence scattering experiments, neutral density filters were added in front of the photomultiplier to keep the detection system linear. Each filter was individually calibrated.

6.4 Results.

A spatial scan across the diameter of HBTX1A was carried out and a profile of the neutral density calculated from the D- α and D- β fluorescence traces. Profiles generated from a transport code, called the AURORA code (Hughes and Post, 1978) were generated to match the calculated profile.

6.4.1 Neutral Density Profile. D- α fluorescence scattering with a resolution of 4 cm was carried out in the plasma region between the inner wall and the plasma centre. D- β fluorescence scattering with a resolution of 4 cm was carried out in the plasma region between the outer wall and the centre. A vertical field of 8.5 mT was then added to move the plasma into a more central position in the vessel and D- β fluorescence scattering carried out again on the outer plasma region with both 4 cm and 2 cm resolution.

A much improved fluorescence enhancement (approximately a factor of 2 larger) was observed upon tuning the laser to the D- β $n=2 \rightarrow 4$ ($\lambda=4860 \text{ \AA}$) transition. This agrees well with the factor of 2 improvement predicted by the CR code. Plate IV v) shows a typical fluorescence trace for both transitions, taken at the same position $r=10.5 \text{ cm}$ either

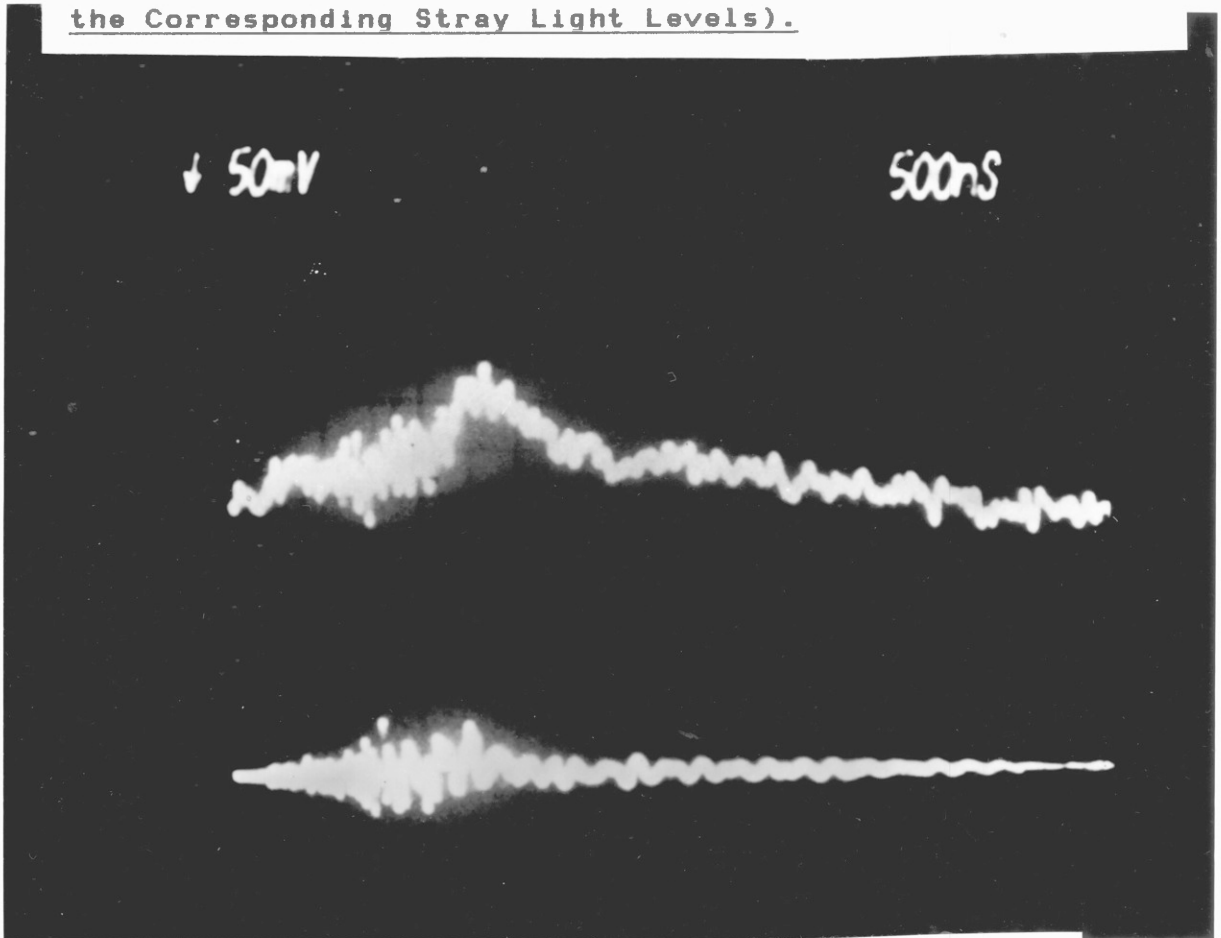
side of the plasma centre. Also shown are the corresponding stray light levels (the laser being fired with no plasma present).

The improvement in moving from the D- α to the D- β transition is clearly seen although the magnitude of the change shown is slightly misleading because the actual neutral density is higher in the outer region of the plasma. Although the two traces are displayed on two different timebases, it can be seen that the D- β trace is approximately 290 ns longer (FWHM) than the D- α trace, due to the longer time duration of the laser at this wavelength.

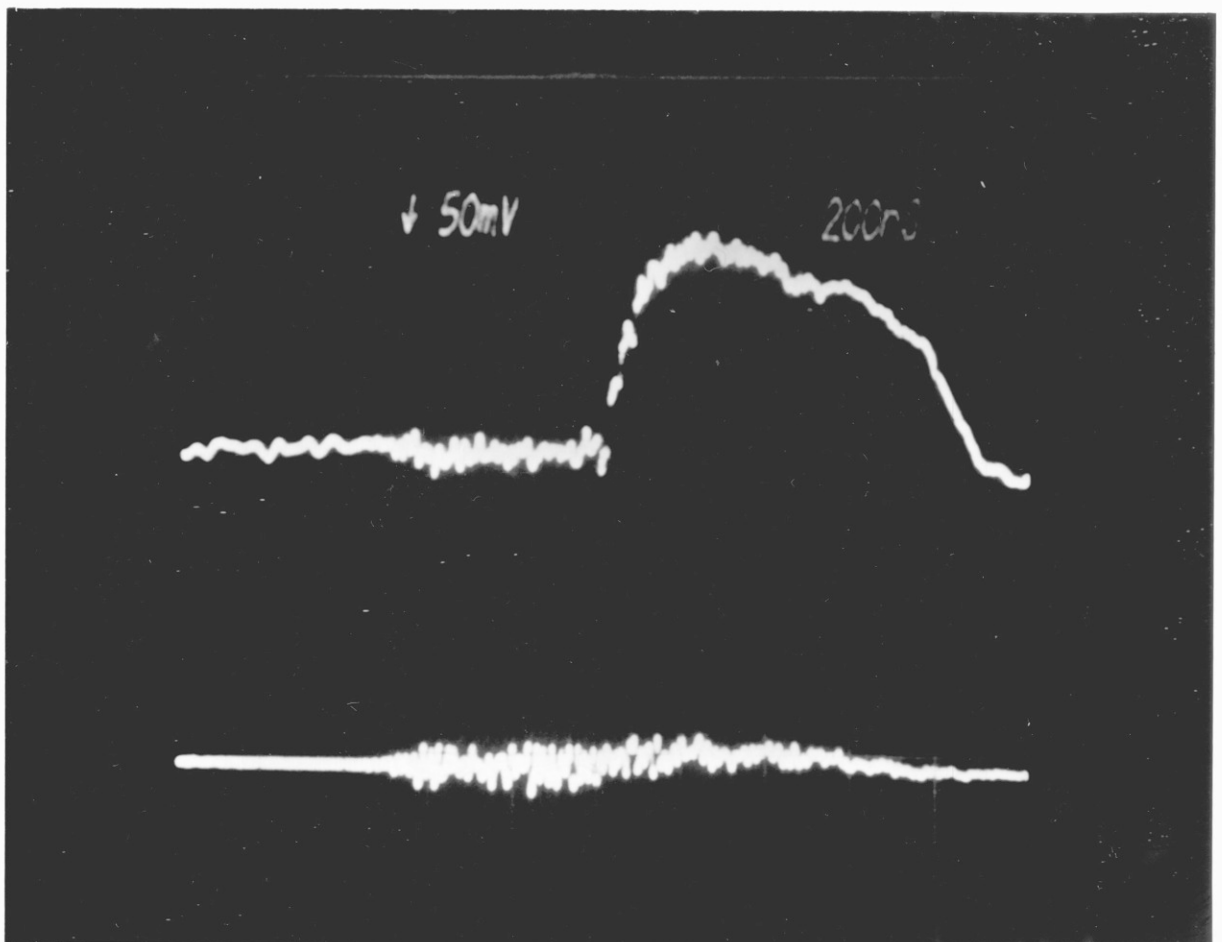
Figure 6.8 shows the experimentally determined neutral density profile in HBTX1A. Each data point was averaged over at least 3 shots for the D- α fluorescence scattering and over at least 10 for the D- β fluorescence scattering results. The error represents one standard deviation from the mean. Under normal operating conditions (no vertical field) the plasma appears to be slightly asymmetric, the neutral hydrogen density being higher towards the outer region.

The effect of the vertical field is to decrease the absolute neutral density, particularly towards the centre of the vessel. For the position nearest the centre ($r=10.5$ cm) the density was observed to decrease from $7.28 \pm 1.79 \times 10^{11} \text{ cm}^{-3}$ to $3.37 \pm 0.74 \times 10^{11} \text{ cm}^{-3}$ for the 4 cm resolution points.

Although spatially resolved neutral density profiles have been obtained using H- α /D- α fluorescence scattering previously (see Section 5.5) this work is believed to be the first measurement of its kind at the D- β wavelength (Evans et al, 1985). The main reason for this is because of the high laser powers required to saturate this transition ($80 \rightarrow 200 \text{ KWcm}^{-2}$ for $n_e = 10^{13} \text{ cm}^{-3}$ and $T_e = 100 \text{ eV}$ as opposed to $4 \rightarrow 10 \text{ KWcm}^{-2}$ for the D- α transition, see



1) Balmer-Alpha Fluorescence Scattering ($r=-10.5$ cm).



11) Balmer-Beta Fluorescence Scattering ($r=10.5$ cm).

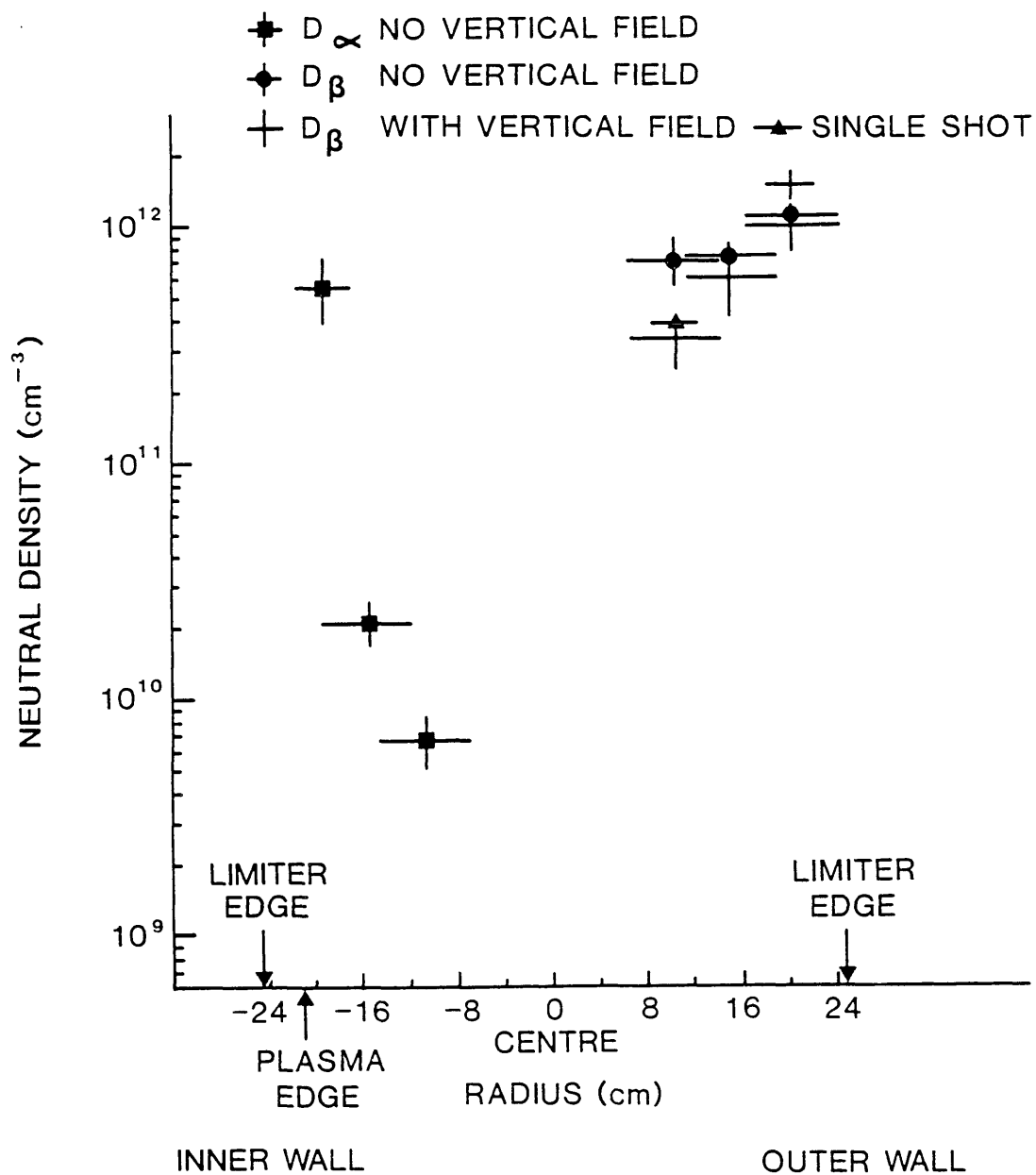


Figure 6.8 Experimentally Determined Neutral Density Profile in HBTX1A from both Balmer-Alpha and -Beta Fluorescence Scattering.

Figure 6.1) coupled with the fact that previous work with the Coumarin dye (Ward, 1983) suggested that it did not perform well when used in flashlamp pumped dye lasers.

6.4.2 Fluorescence Enhancement. Table 6.2 shows the fluorescence enhancements measured (n_3^L/n_3^0 , n_4^L/n_4^0) for each position investigated plus the theoretical values for the given plasma conditions from the CR code. The positions are in relation to the plasma centre ($r=0$ cm) and are negative for the region between the centre and the inner wall. Each observed enhancement has been corrected by the appropriate geometric factor, and the error quoted represents one standard deviation. The D- α fluorescence enhancements measured agree with the expected values to within experimental error. The D- β fluorescence enhancements are also scattered about the predicted values. An enhancement larger than theory by a factor of approximately 2 was obtained for the position at $r=20.5$ cm with the vertical field and 2 cm resolution. However, only two experimental shots were taken at this position and the background level was observed to be lower than the previous shots by a factor of 2. The plasma may have broken down early in this case and these experimental points may correspond to a later time in the discharge. This data has been included in Figure 6.8, but the error bars have been calculated from the fluctuation in the plasma background.

6.4.3 Local Electron Density Measurement. An attempt was made to extend the diagnostic applications of the fluorescence scattering method. As discussed by Gohil (1982), if the laser is cut off sharply in time whilst still saturating the transition, the fluorescence signal will decay in a time given by the main rate out of the upper level, which is the electron collision rate $n_e C_{j1}$. As the collisional coefficient is well known (there is no ambiguity as regards the exact rate coefficients at temperatures well above threshold) then the local electron density may be calculated from the measured

TABLE 6.2: The Fluorescence Enhancements
Observed in the HBTX1A Plasma Device.

Radius (cm)	T_e (eV)	n_e (10^{14} cm^{-3})	Measured Enhancement	Theory (CRcode)	Note
-19.0	18.3	0.55	2.4 ± 1.0	2.6	(i)
-15.5	63.9	1.92	2.5 ± 0.7	2.7	(i)
-10.5	87.5	2.63	3.8 ± 1.0	2.8	(i)
10.5	99.7	2.99	8.0 ± 3.4	5.3	(ii)
15.5	96.1	2.88	6.6 ± 2.5	5.4	(ii)
20.5	74.8	2.24	5.0 ± 1.4	5.5	(ii)
10.5	99.5	2.98	3.3 ± 0.8	5.3	(iii)
15.5	95.5	2.86	6.3 ± 1.6	5.4	(iii)
20.5	77.3	2.32	6.6 ± 2.6	5.5	(iii)
10.5	99.5	2.98	5.6 ± 0.6	5.3	(iv)
20.5	77.3	2.32	13.1 ± 1.1	5.5	(v)

(i) D- α fluorescence scattering, $B_v=0$, 4 cm resolution.

(ii) D- β fluorescence scattering, $B_v=0$, 4 cm resolution.

(iii) D- β fluorescence scattering, $B_v=8.5$ mT, 4 cm resolution.

(iv) D- β fluorescence scattering, $B_v=8.5$ mT, 2 cm resolution, single shot error estimated from plasma background variation.

(v) Same as (iv) above, but with 2 experimental shots only, error estimated from plasma background variation.

decay time. The importance of knowing the exact form of the electron density profile in this type of plasma is discussed in the next section (Section 6.5.2). For the range of densities expected in HBTX1A ($1 \rightarrow 3 \times 10^{14} \text{ cm}^{-3}$) the D- β transition gives the largest decay time variation plus decay times that are not too small to measure ($2 \rightarrow 10 \text{ ns}$). The D- α fluorescence decay for these plasma conditions although measurable ($7 \rightarrow 9 \text{ ns}$ Gohil, 1982), does not give a wide enough range for a variation in the electron density to be observed and the Balmer-gamma fluorescence signal decays in times of less than 2 ns (Gohil, 1982) which would be experimentally very difficult if not impossible to measure.

This experiment was attempted after the neutral profile had been obtained. By careful alignment of the Pockel cell, more than 90% extinction was obtained for the falling edge of the laser pulse in a time $\leq 2.5 \text{ ns}$, measured with an ITL fast photodiode (risetime at $3 \text{ KV} \leq 100 \text{ ps}$) and displayed on a 7104 Tektronix oscilloscope (7B15 plugin with a quoted risetime of 300 ps) connected by 50 cm of a special low attenuation cable supplied by Tektronix. Unfortunately, the fall time of the photomultiplier used was found to be too long (the fall time of the laser itself was measured to be approximately 6 ns with this photomultiplier) and the experiment could not be completed because the fast 7104 oscilloscope (which was on loan) had to be returned to Tektronix.

6.5 Discussion.

6.5.1 The AURORA Code. For comparison with the results, theoretical neutral density profiles were generated using a transport code called the AURORA code (Hughes and Post, 1978).

The AURORA code is a Monte Carlo algorithm which calculates neutral gas transport in cylindrical plasmas,

assuming a monoenergetic isotropic source. This assumption is widely used in transport codes. A particle is launched at the boundary and its path tracked until it undergoes ionization (in which case the particle is considered lost) or charge exchange with protons (in which case the particle is relaunched with a velocity chosen from a Maxwellian distribution characterised by the local ion temperature). The electron ionization rates are obtained from the polynomial fit given by Freeman and Jones (1974) and the charge exchange reaction rates from formulae in Riviere (1971).

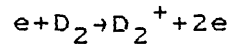
1) Temperature and Density Profiles. The AURORA code assumes a symmetric plasma with ion and electron temperature and electron density radial profiles of the form:

$$T_e(r) = T_e(0)(1 - (r/a)^N)^{M+Pd} \quad (6-7)$$

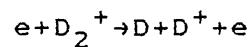
where $T_e(r)$ is the electron temperature corresponding to the radius r in the plasma, $T_e(0)$ is the peak temperature in the centre ($r=0$) and a is the plasma radius ($a=24.5$ cm for HBTX1A as defined by the limiter edge). Pd is the pedestal and is equal to the value of the edge ($r=a$) electron temperature. The electron density and the ion temperature profiles are of the same form.

ii) Origin of the Neutral Particles. The AURORA code assumes a monoenergetic source of neutral particles. Of the neutrals which bombard the walls, the atoms are more likely to stick than the molecules because of the much higher binding energy of the atoms (Vijh and Belanger, 1973). It seems likely that since the walls are loaded with molecular deuterium before the plasma discharge begins, this is released from the walls due to neutral and electron bombardment. The molecular deuterium may then undergo several reactions to produce atomic deuterium. For the range of temperatures expected at the edge of HBTX1A (10-100 eV) the process with the largest

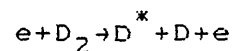
rate per molecule (by an order of magnitude, from Kaplan and Bengtson (1981)) is direct molecular ionization:



the D_2^+ molecule is then most likely to produce atomic deuterium via molecular ion dissociation where the neutrals produced have energies of a few eV (see Section 5.3):



The other main process involving the formation of neutral atoms is electron impact dissociation of D_2 (Kaplan and Bengtson, 1981):



This reaction produces atoms belonging to one of two distinct energy groups. Slow D atoms with energies of approximately 0.3 eV are formed in states with $n \geq 2$ by electrons with energies ≥ 15 eV and fast D atoms having energies in the range 2-8 eV are formed by electrons with energies ≥ 28 eV (Levanthal, 1967).

Most of the reactions involving the formation of neutral atoms have thresholds exceeding 10 eV, and so the rate coefficients are extremely sensitive to the lower temperatures near the plasma edge. Accurate estimates of the edge temperature is essential for modelling of the plasma formation in this region.

6.5.2 Radial Electron Temperature and Density Profiles.

It is important to know the exact temperature and density profiles for the generation of neutral profiles with transport codes. For the fluorescence results themselves the output from the CR code for $T_e > 30$ eV is only very weakly dependent on the profiles chosen, as discussed in Section 5.3.2. Results from the AURORA code are sensitive

to the type of profile chosen, in particular to the form of the electron density profile.

For the results presented so far, the electron temperature and density profiles were chosen from Malacarne (1984) this being the most recent and reliable data available. His results correlate plasma emission (mainly due to CV observed with xray diodes) with magnetic edge coil fields and shows that there is a rapid radial propagation of thermal energy in the centre followed by a slower diffusion of the excess temperature profile in the edge region. This suggests a flat-topped electron temperature distribution in the centre, falling off rapidly at the plasma edge. Figure 6.9 shows the radial line of sight xray emission (Abel inverted) for two HBTX1A shots taken from Malacarne. Note that the discharge is shifted away from the inner wall by 3 cm for normal operating conditions (i.e no vertical field). This shift is seen throughout the entire discharge time. It is predicted by toroidal equilibrium calculations (Mukhovatov and Shafranov, 1971) and has been confirmed on HBTX1A by electrical diagnostics and magnetic probe measurements (Radcliffe, 1984). If an external vertical magnetic field is applied the shift may be reduced or eliminated. The effect of such a vertical field ($B_V=8.5$ mT) on the xray emission is also shown in Figure 6.9. The intensity distribution is seen to become virtually symmetric around the central position.

It was decided to use Malacarne's data as a guide to the actual shape of both the electron and ion temperatures and the electron density, with a peak electron temperature of 100 eV (from bolometer results, P.G. Carolan, private communication) and peak density of $3 \times 10^{13} \text{ cm}^{-3}$ (CO_2 laser interferometer results, D.E. Evans private communication). Further, since the temperature profile obtained from Malacarne was asymmetrical about the centre point for normal HBTX1A operating conditions, the two halves of the plasma were considered

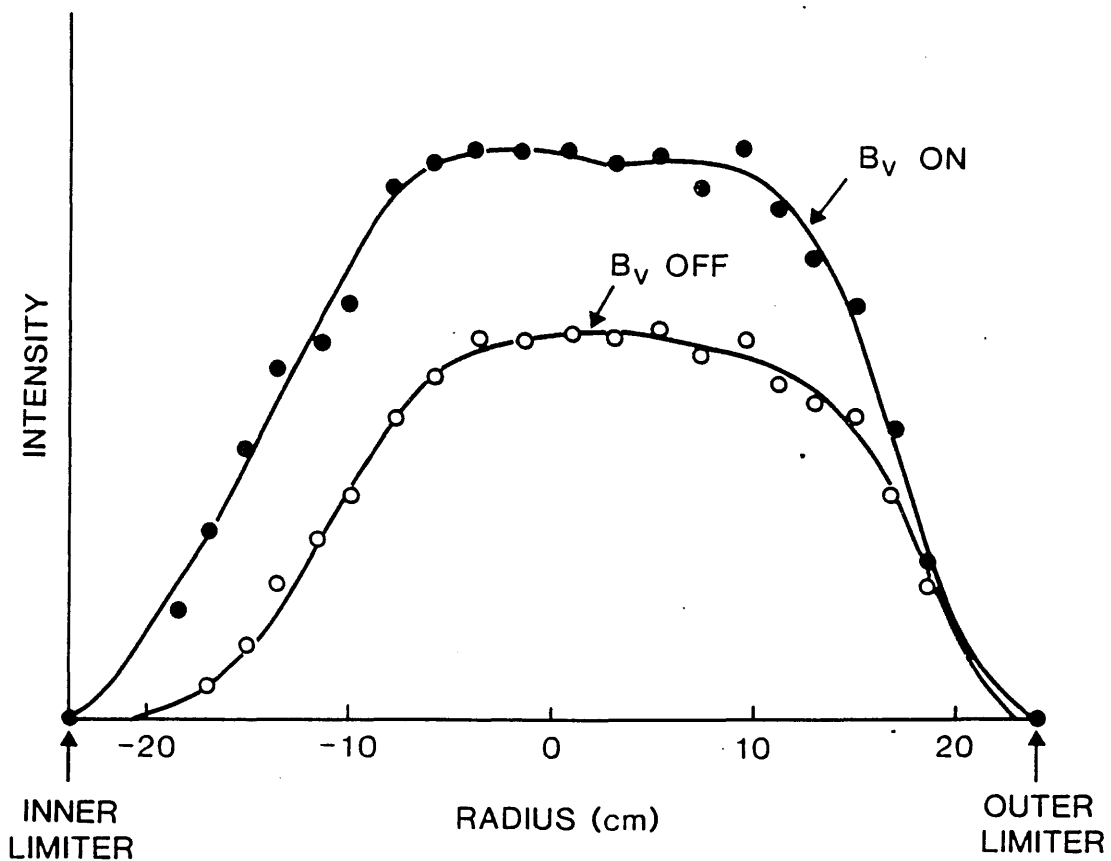


Figure 6.9 Radial Line of Sight X-Ray Emission from HBTX1A.

This shows the Abel inverted xray emission from HBTX1A from Malacarne (1984) for a typical shot (no vertical field) and also a shot with the vertical field applied. Note the more symmetrical profile when B_v is applied.

separately. Prior to discussing the results of the AURORA code some details pertaining to its use must be discussed.

1) HBTX1A Geometry. For normal operating conditions (with no vertical field) the plasma is asymmetrically distributed, lying 3 cm away from the inner limiter. To use the AURORA code the coordinate system must be redefined so that symmetry is maintained. The required transformation is to define a new radius:

$$r' = r - 1.5 \text{ cm}$$

with the radius of the plasma $a = 23$ cm. Figure 6.10 shows the geometry of HBTX1A to make this transformation clear.

To clarify the discussion of the various regions of the plasma, the following definitions describe those parts of the plasma studied experimentally:

1. INNER REGION, with no vertical field applied. This was considered to be the plasma from radius $r' = 0$ to -23 cm under usual HBTX1A operating conditions. The plasma "edge" was 3 cm from the inside limiter. D- α fluorescence measurements were carried out in this region.

2. OUTER REGION, with no vertical field applied. This was considered to be the plasma region from $r' = 0$ to $+23$ cm, again under usual HBTX1A operating conditions. The plasma "edge" was in contact with the outside limiter at $r' = 23$ cm. D- β fluorescence measurements with 4 cm spatial resolution were carried out in this region.

3. OUTER REGION, with a vertical field applied. This region defines the plasma between $r = 0$ to 24.5 cm (N.B. $r = 0$ is the geometric centre of the plasma) with a vertical field B_v of 8.5 mT applied. D- β fluorescence measurements were performed in this region with both 2

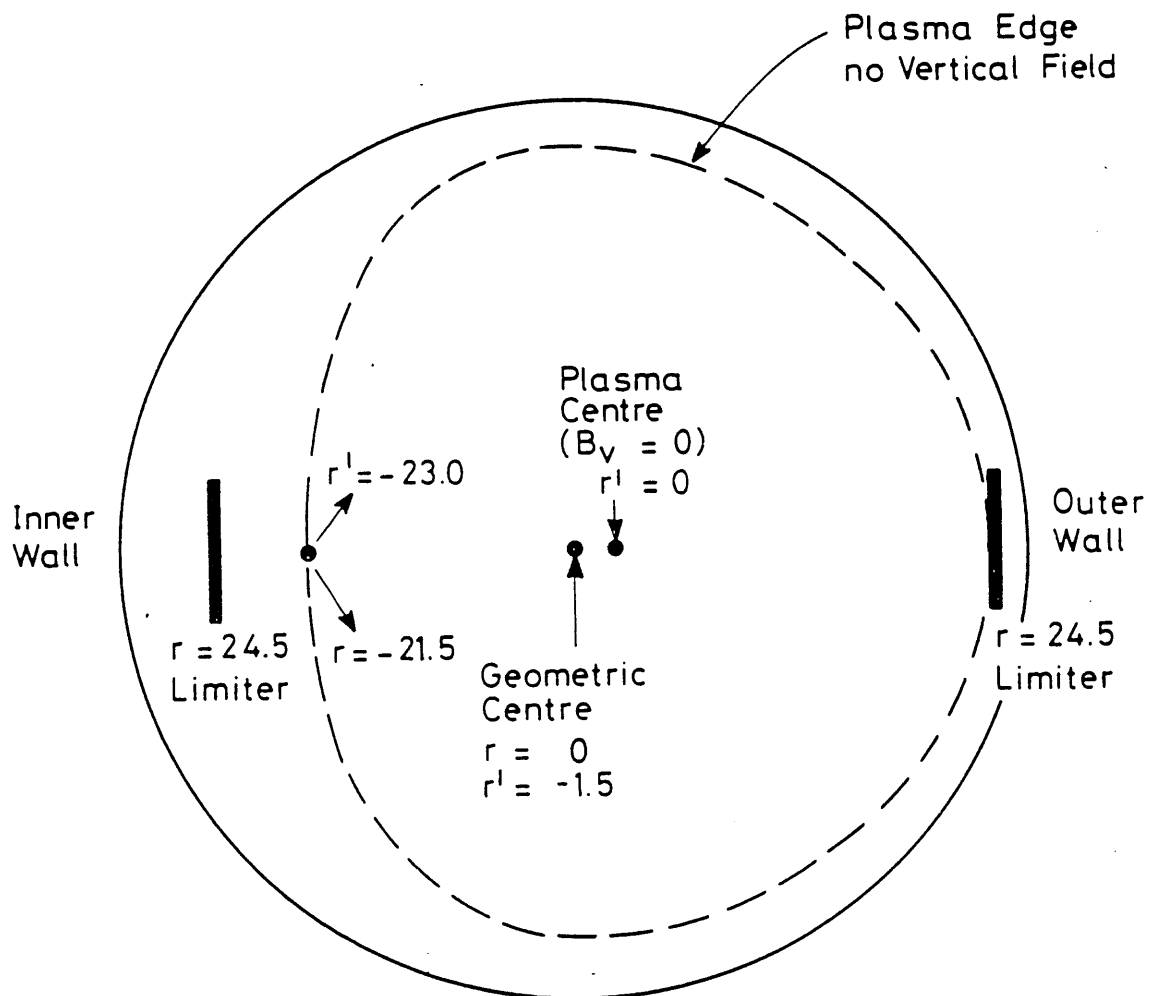


Figure 6.10 The Geometry of the HBTX1A Cross Section.

This shows the position of the plasma within the HBTX1A vessel when no vertical field is present (i.e. normal operating conditions). Note that the plasma centre is not coincident with the geometric centre of the vessel, the plasma edge being 3 cm from the inner wall limiter.

and 4 cm spatial resolution.

ii) Parabolic Temperature and Density Distribution. When no other data is available this is the distribution usually chosen to describe the radial variation in the temperature and density. For this distribution, $N=2$ and $M=1$ in Equation 6-7.

For a parabolic distribution, with the radius redefined as above and the pedestal level $P_d=0$:

$$T_e(r') = T_e(0)(1 - |r'/a|^2) \quad (6-8)$$

where all the symbols have been previously defined.

iii) Best Fit to Existing Data Distribution. The following analysis was performed to find the N and M from Equation 6-5 which would give the best fit to the available data from Malacarne (1984). Taking Equation 6-5 and redefining the new radius r' as above gives, for $P_d=0$:

$$T_e(r')/T_e(0) = (1 - |r'/a|^N)^M \quad (6-9)$$

Now the binomial expansion is of the form:

$$(1+x)^\alpha = 1 - \alpha x + \alpha(\alpha-1)x^2/2! - \alpha(\alpha-1)(\alpha-2)x^3/3! + \dots$$

Expanding Equation 6-7 as a binomial series, and replacing $T_e(r')/T_e(0)$ by the normalised temperature variable $T_e'(r')$ gives:

$$\begin{aligned} T_e'(r') = & 1 - M|r'/a|^N + \frac{M(M-1)}{2!}|r'/a|^{2N} \\ & - \frac{M(M-1)(M-2)}{3!}|r'/a|^{3N} + \dots \end{aligned} \quad (6-10)$$

This is valid for $(r'/a) < 1$, which is always satisfied since $r'=a$ corresponds to the edge of the plasma. For $N > 1$, this can be approximated to:

$$T_e'(r') \approx 1 - M|r'/a|^N$$

Rearranging and taking logs gives:

$$\ln(1-T_e'(r')) = N \ln|r'/a| + \ln M \quad (6-11)$$

This is equivalent to a straight line of the form $y=mx+c$ where $y=\ln(1-T_e(r'))$, $x=\ln(r'/a)$, the gradient $m=N$ and the intercept $c=\ln(M)$.

For each region of the plasma (as defined above), a linear regression was performed to determine the coefficients N and M which would give the best fit between the experimental data and Equation 6-7. Those points where, $T_e(r')=T_e(0)$ were excluded to prevent a singularity in Equation 6-11 at $\ln(0)$. Table 6.3 gives the values of the coefficients so obtained for the inner and outer regions of the plasma under ordinary HBTX1A operating conditions, and for the outer region with the vertical field present.

6.5.3 Comparison Of Fluorescence Results and AURORA Output.

1) AURORA Code. The effect of the different radial temperature and density profiles on the output from the AURORA code and the calculated neutral densities from the fluorescence results was investigated. Figure 6.11 shows neutral density profiles generated from the AURORA code for a parabolic temperature and density distribution for a range of neutral input energies. The incoming flux of neutrals in each case has been normalised so that the external density is the same (10^{11} cm^{-3}). The typical error bars shown are generated from the AURORA code, increasing towards the centre to a typical error of 15%. The neutral profiles are more sensitive to the lower input energies ($<5 \text{ eV}$).

Figure 6.12 shows the neutral profiles generated for the best fit distributions for each region of HBTX1A studied.

TABLE 6.3: The Exponential Coefficients N and M as
Deduced From a Best-Fit of Available Data to Equation
6-7 for Radial Temperature and Density Profiles.

Region	Plasma radius (cm)	N	M	Correl. Coeff.
Inner ($B_V=0$)	23	3.27	1.05	0.9946
Outer ($B_V=0$)	23	5.95	0.75	0.9946
Outer ($B_V=8.5$ mT)	24.5	5.45	0.54	0.9954

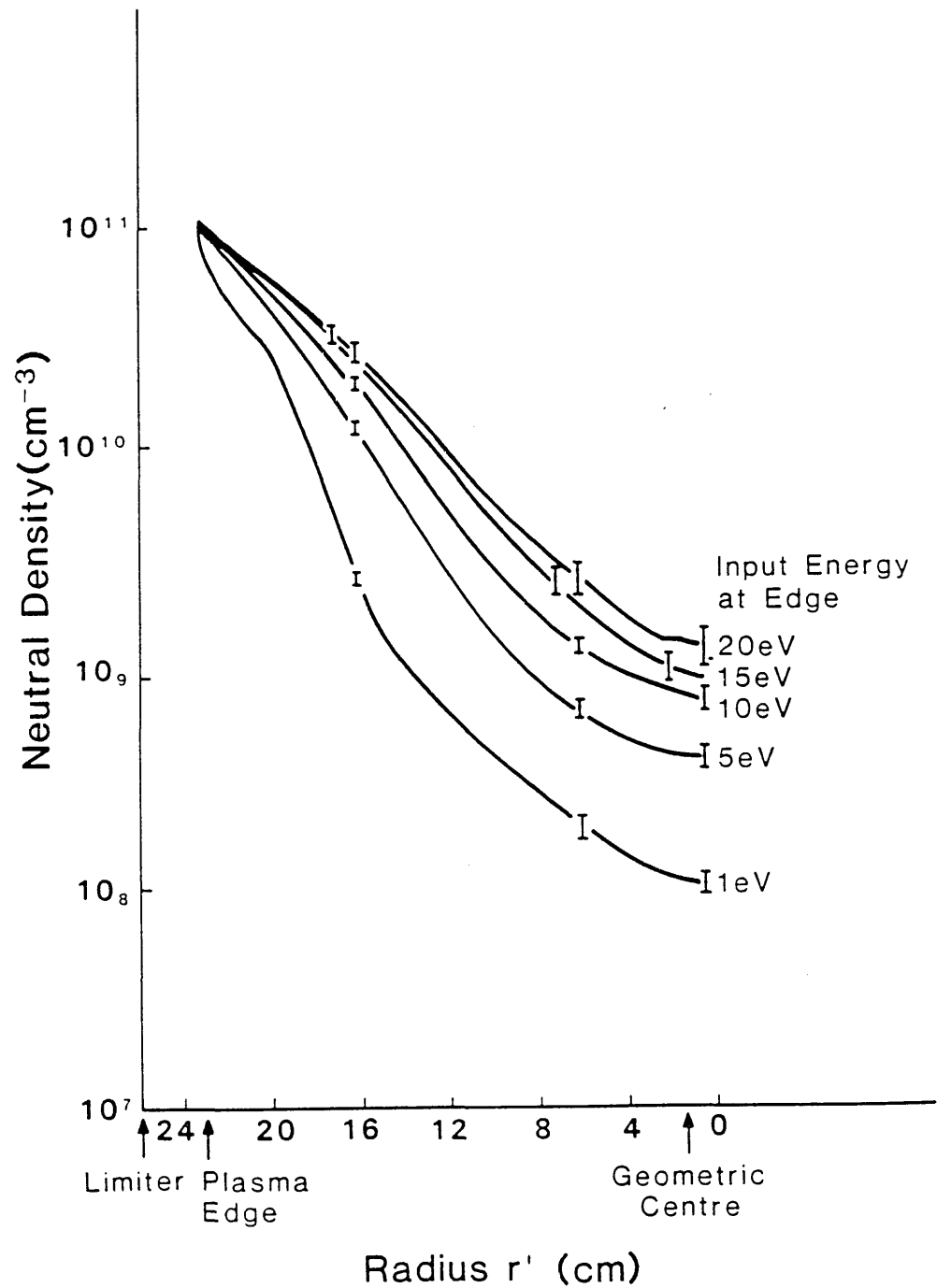


Figure 6.11 1) Neutral Density Profiles Generated from the AURORA Code for a Parabolic Temperature and Density Distribution for the Inner Region of HBTX1A.

This set of profiles was generated for a range of neutral energies at the plasma edge. The edge density has been set to be 10^{11} cm^{-3} in all cases and the radius $a=23 \text{ cm}$ (for normal operating conditions) and pedestal $P_d=0$.

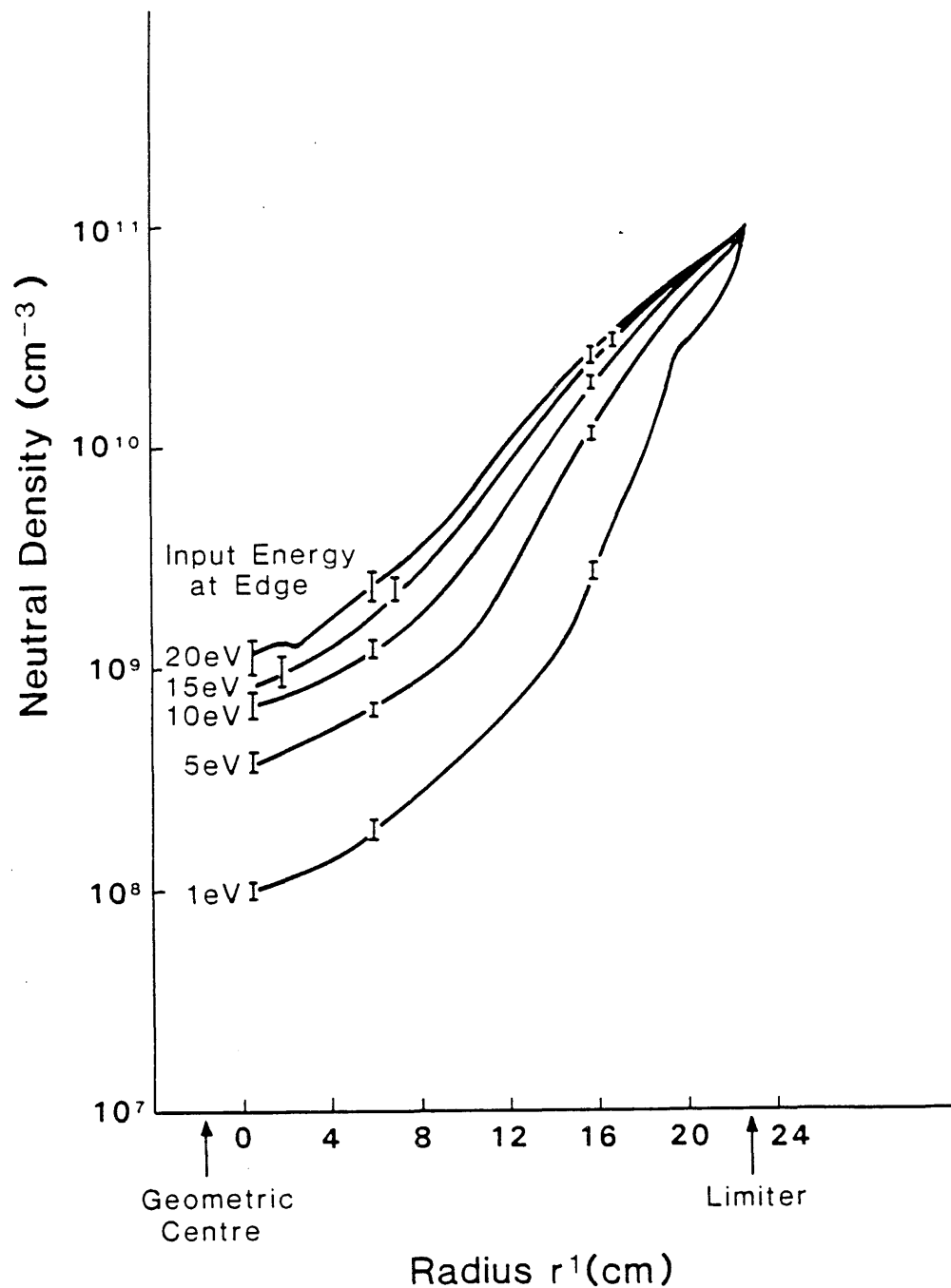


Figure 6.11 11) Neutral Density Profiles Generated from the AURORA Code for a Parabolic Temperature and Density Distribution for the Outer Region of HBTX1A.

This set of profiles was generated for a range of neutral energies at the plasma edge. The edge density has been set to be 10^{11} cm^{-3} in all cases and the radius $a=23 \text{ cm}$ (for normal operating conditions) and pedestal $P_d=0$.

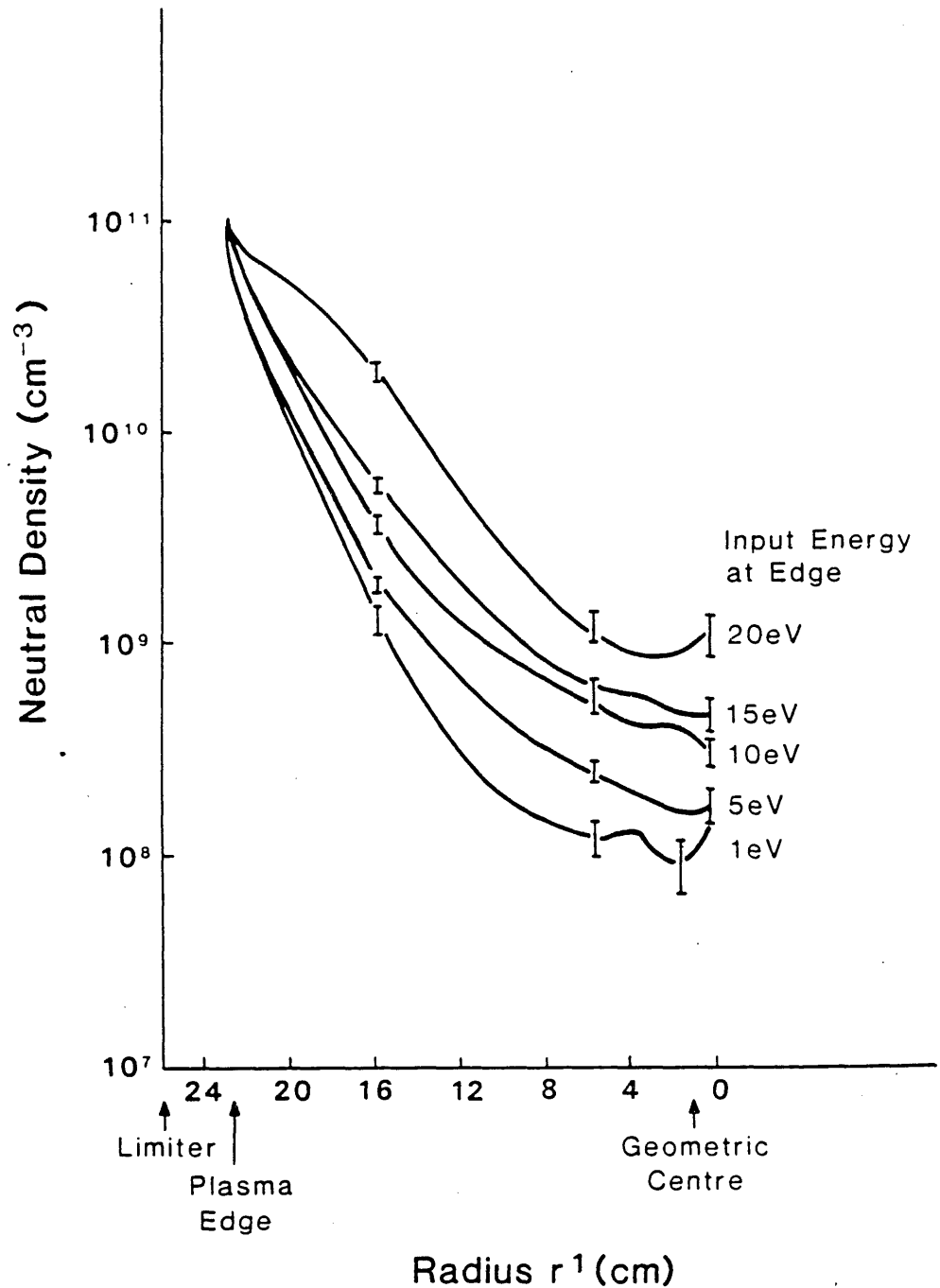


Figure 6.12 1) Neutral Density Profiles Generated from the AURORA Code for a Best Fit Temperature and Density Distribution for the Inner Region of HBTX1A.

This set of profiles was generated for a range of neutral energies at the plasma edge. The edge density has been set to be 10^{11} cm^{-3} in all cases and the radius $a=23 \text{ cm}$ (for normal operating conditions) and pedestal $P_d=0$. The radial electron temperature and density profiles used in the code were from a fit to the data of Malacarne (1984).

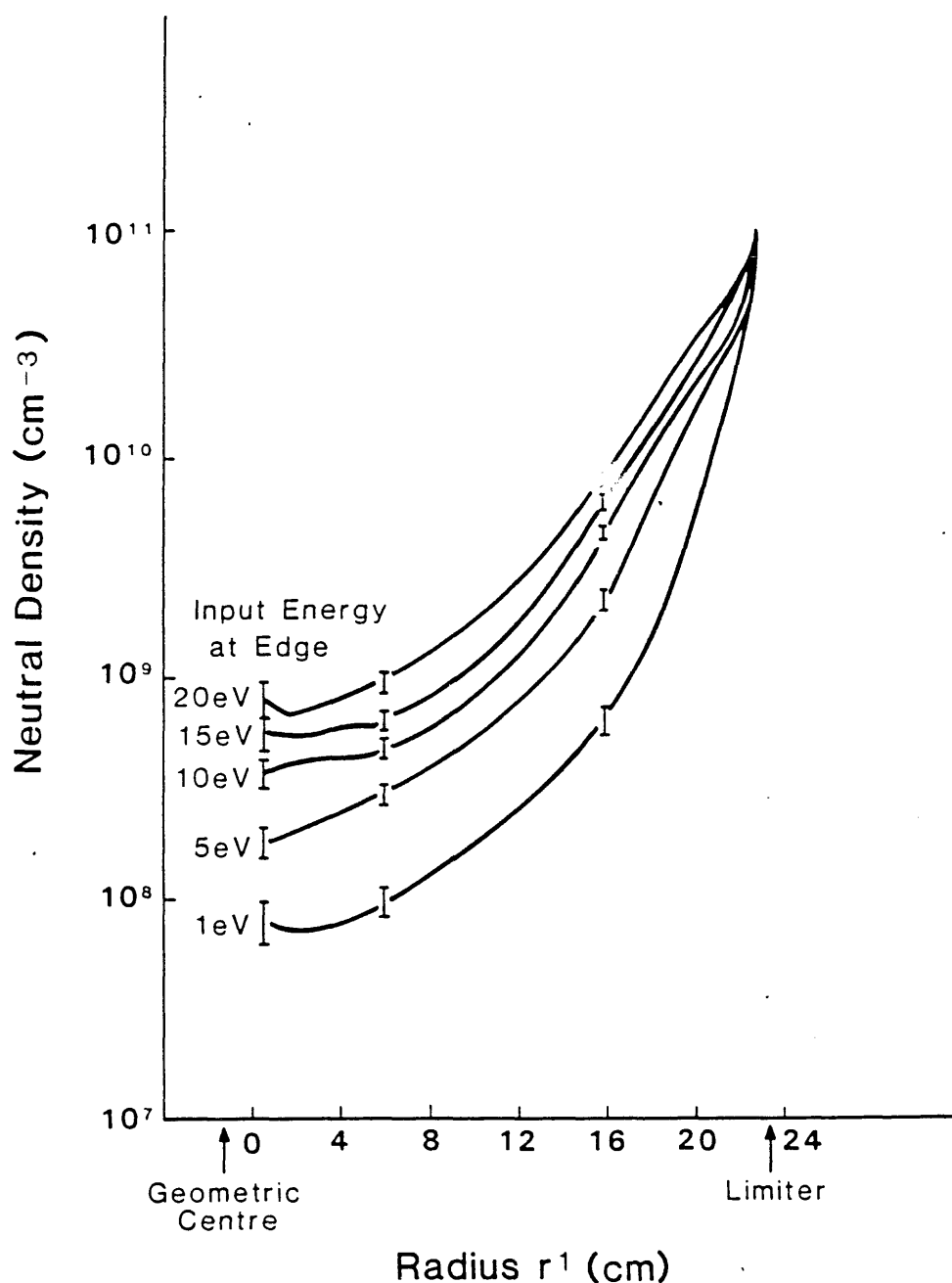


Figure 6.12 11) Neutral Density Profiles Generated from the AURORA Code for a Best Fit Temperature and Density Distribution for the Outer Region of HBTX1A.

This set of profiles was generated for a range of neutral energies at the plasma edge. The edge density has been set to be 10^{11} cm^{-3} in all cases and the radius $a=23 \text{ cm}$ (for normal operating conditions) and pedestal $P_d=0$. The radial electron temperature and density profiles used in the code were from a fit to the data of Malacarne (1984).

Again, the profiles appear to be particularly sensitive to the lower energy input neutrals (<5 eV). Comparing Figure 6.11 with Figure 6.12 for a given input neutral energy, the profiles generated by the best fit temperature and density distributions fall off much more rapidly towards the centre than those generated from a parabolic distribution. This is because the best fit distributions have a flat-topped shaped profile in the plasma centre. Being hotter and more dense than in the case of the parabolic distribution, the neutral particles are more likely to undergo ionization and/or charge exchange.

11) Fluorescence Results. The experimentally determined neutral density profile shown in Figure 6.8 was obtained using the radial electron density and temperature profiles of Malacarne (1984). Use of a parabolic density and temperature profile in the CR code gives higher values of the neutral density. For the $r=\pm 10.5$ cm point (the most central position investigated) the neutral density is less than 15% higher than when the best fit profile is used. For the position at ± 15.5 cm the neutral density is typically 30% larger and for the position nearest the wall ($r=+20.5/-19.0$) is typically a factor a factor of 2 larger.

6.5.4 Analysis of Data. To compare the experimentally determined neutral density profiles with the output from the AURORA code, the following analysis was performed.

The neutral density profiles generated by the AURORA code for given temperature and density distributions were assumed to have a radial distribution of the form:

$$n_1(r') = n_1(0)e^{r'/R} \quad (6-12)$$

where $n_1(r')$ is the neutral density at radius r' , $n_1(0)$ is the central ($r'=0$) neutral density and R is a constant.

$$\ln(n_1(r')) = (r'/R) + \ln(n_1(0)) \quad (6-13)$$

A linear regression of the form $y=mx+c$ where $y=\ln(n_1(r'))$, $x=r'$, the gradient $m=1/R$ and the constant $c=\ln(n_1(0))$ was performed to find the values of the constants R and $n_1(0)$ which obtained the best fit of the AURORA-generated profiles to Equation (6-12) for the values of r' over which experimental data was available. Table 6.4 shows the values of these coefficients deduced from AURORA output for the three plasma regions investigated. The value of the constant $n_1(0)$ obtained from the linear regression is arbitrary, depending on the edge density $n_1(a)$, which is fixed to have the value 10^{11} cm^{-3} within the code. The constant R is independent of the edge density chosen.

The number of data points used for each regression presented in Table 6.4 was 13. For 13 data points, the correlation coefficient which gives a 99% certainty that the data is not due to chance is 0.684 (Bevington, 1969). All the correlation coefficients listed in Table 6.4 are greater than this, so the assumption that the AURORA profiles (over the range which data was taken) fit an exponential is good.

Taking the experimental values of $n_1(r')$ as measured from the fluorescence traces, a linear regression was performed on each set, fitting the data to an exponential of the same form as Equation (6-12). The result of this is shown in Table 6.5 for the neutral density calculated for both the parabolic and best fit density and temperature profiles. For the inner region, the parabolic and best fit correlation coefficients suggest that the fits are more than 80% significant. For the outer region with no vertical field present, the parabolic fit is more than 80% significant and the best fit more than 50% significant, so the data does not fit to an exponential as well as the inner region data. For the outer region

TABLE 6.4: The Coefficients Produced by Fitting Output from the AURORA Code to the Exponential Distribution given in Equation 6-11.

REGION	VERT.FIELD (ON/OFF)	DISTRB.	ENERGY OF EDGE INPUT NEUTRALS	R (cm)	CORREL. COEFF.
Inner/ Outer	Off	Parabol.	15 eV	4.39	0.9868
			10	3.71	0.9947
			5	3.00	0.9972
			1	2.34	0.9893
Inner	Off	Best fit	15 eV	3.42	0.9974
			10	3.25	0.9560
			5	2.57	0.9947
			1	1.05	0.9850
Outer	Off	Best fit	15 eV	3.06	0.9937
			10	2.60	0.9906
			5	2.59	0.9653
			1	2.13	0.9270
Outer	On	Parabol.	15 eV	4.39	0.9868
			10	3.71	0.9947
			5	3.00	0.9972
			1	2.34	0.9893
Outer	On	Best fit	15 eV	3.34	0.9943
			10	3.24	0.9976
			5	2.75	0.9958
			1	2.35	0.9340

TABLE 6.5: Coefficients from the Best Fit of Experimental Data Points to an EXponential Given in Equation (6-12).

REGION	VERT.FIELD (ON/OFF)	DISTRIB.	R (cm)	N (cm ⁻³)	CORREL. COEFF.
Inner	Off	Parabol.	1.79	1.5E06	0.9299
Inner	Off	Best Fit	2.00	2.4E07	0.9280
Outer	Off	Parabol.	9.43	2.8E11	0.9378
Outer	Off	Best Fit	23.26	4.6E11	0.8859
Outer	On	Parabol.	5.46	6.3E10	0.9745
Outer	On	Best Fit	8.77	1.1E11	0.9949

(N.B. 1Ex=10^x)

TABLE 6.6: Estimate of the Energy of the Edge Neutrals from a Comparison Between the Data and Output from the AURORA Code.

REGION	VERT.FIELD (ON/OFF)	DISTRIB.	ENERGY OF EDGE NEUTRALS (eV)
Inner	Off	Parabol.	0.0
Inner	Off	Best fit	5.5
Outer	Off	Parabol.	>100.0
Outer	Off	Best fit	>100.0
Outer	On	Parabol.	>100.0
Outer	On	Best fit	>100.0

with the vertical field present, the parabolic fit gives approximately 80% significance and the best fit 90% significance.

By comparing the constants R obtained in Table 6.5 with those from the AURORA code in Table 6.4, an estimate of the energy of the edge neutrals was made (note that this constant is independent of the edge density chosen because this represents the slope of the data). The results are tabulated for each region of the plasma in Table 6.6. At the outer region, with or without a vertical field present, the edge neutral energy predicted by the AURORA code is not realistic.

6.6 Conclusions.

A spatially resolved neutral density profile was obtained across the radius of HBTX1A. Assumptions had to be made concerning the exact temperature and density profiles and any conclusions drawn from the preceding computational analysis must be treated with caution because of this.

6.6.1 Neutral Density Profile. The application of the AURORA code in this case does not appear to be valid as regards predicting the energy of the edge neutrals, particularly in the outer plasma region, since the edge neutral energy would have to be greater than 100 eV (and this is the peak electron temperature in the plasma centre). Only one positive result was obtained. If the radial density and temperature profiles are fitted to the best data available (Malacarne, (1984), who predicted a flat-topped distribution) the input energy of the edge neutrals from the inner wall is found to be 5.5 eV. This energy would be consistent with the formation of neutrals from electron impact ionization of molecular deuterium.

It should be noted that the code does not include the effect of charge exchange with impurities. If this process was included, the absolute value of the neutral

density generated at a given position in the plasma would be decreased although the slope of the neutral profile (the constant R in Equation (6-12)) may not change. For the generated profiles to fit the experimental data, charge exchange would have to occur predominantly in the outer regions of the plasma, near to the wall and inclusion of charge exchange with impurities into the AURORA code may well produce a better result as regards the comparison with theory. This result would also suggest that the impurities present were produced by sputtering, and this needs further investigation.

Addition of a vertical field caused the observed neutral density to decrease, this effect being more pronounced for the most central position in the plasma. This result confirmed previous experimental work which showed that addition of such a field caused the plasma to shift to a more central position in the vessel, away from the outer wall (Radcliffe, 1984). This is also consistent with the premise that the origin of the neutrals in the plasma is primarily from plasma-wall interactions, such as the collisional dissociation of D_2^+ which itself originates from the ionization of molecular deuterium at the wall.

6.6.2 Fluorescence Enhancement. Gohil et al (1983) measured the $D-\alpha$ fluorescence enhancement in HBTX1A and found that it was a factor of 4 larger than theory. So the values of the fluorescence enhancements for this set of experiments were measured for each position in the plasma and compared to the theoretical values from the CR code. Both the $D-\alpha$ and $D-\beta$ enhancements observed agreed to within experimental error with the predicted values, with the exception of one experimental point where only two shots were taken.

Chapter Seven

FLUORESCENCE SCATTERING AS A DIAGNOSTIC FOR IMPURITY CONCENTRATIONS.

In this chapter, the application of fluorescence scattering to the determination of impurity densities in a fusion plasma is discussed.

A feasibility study is presented, which considers the forbidden $^3P_1 - ^3P_2$ ($\lambda=3834.6$ Å) and $^3P_0 - ^3P_1$ ($\lambda=3370.6$ Å) transitions of a particular impurity ion Ti XVII. The study shows that the major experimental obstacle to be overcome is the high laser power required to saturate the chosen transition at the appropriate wavelength.

7.1 The Importance of Impurity Ion Concentration to Fusion Plasmas.

Fusion device plasmas contain impurities such as carbon, oxygen, and certain transition metals derived from the walls and limiters, mostly due to sputtering from high energy neutrals. Under typical plasma conditions, these impurities mainly radiate in the UV and X-ray regions of the spectrum. Impurity radiation is of particular importance as a power loss mechanism, radiating a large fraction of the heat input away (see Section 5.3). Studies on the DITE tokamak at the UKAEA Culham Laboratory (Bickerton, 1977) for example, have indicated that heavy metals in particular can account for 97% of the radiated power and 76% of the mean ion charge excess Z_{eff} .

As previously discussed (Section 5.3) when the neutral hydrogen density exceeds 10^{-5} of the electron density (Peacock, 1980) charge exchange may become the dominant recombination process for impurities:



This process is important when considering neutral beam injection, since it lowers the effective ion charge Z_{eff} and increases the radiation losses in the core plasma due to recombination of fully stripped ions (Suchewer et al, 1980). Charge exchange with neutral hydrogen can also impair beam transport and therefore increase radiation losses at the edge of the plasma. The effect of this radiation loss is greatest when the dominant impurity ion species has one or two electrons (Peacock, 1980), and so can be most important in localised regions.

As discussed in Section 5.3, the centre of tokamak plasmas is often "thick" to emerging neutrals, and so measurement of highly ionized impurity ion concentrations may be exploited as a diagnostic for these regions.

To estimate impurity ion concentration, several methods can be used. Observation of spectral line intensities is perhaps the most common, (Afrosimov et al (1978), Suchewer et al (1980)) although this method has limited spatial resolution. Another technique, described in Section 5.6 was utilized by Afrosimov et al (1978) which involved using a modulated neutral hydrogen beam to perturb the plasma. The concentration of fully ionized carbon was calculated by observing the change in the emission of the C VI line (due to charge exchange processes).

A full description of the advantages of the fluorescence scattering technique compared to other methods is presented in Section 5.4. This method as applied to an impurity ion transition has much the same advantages as when applied to neutral density determination (i.e. well defined spatial resolution).

7.2 Feasibility Study.

It was decided to investigate the feasibility of using laser induced fluorescence as a technique to measure the number density of specific highly ionized species. These species often possess dipole forbidden transitions, in their ground configurations, of low A values but at wavelengths accessible to lasers.

7.2.1 Introduction. The particular species chosen for the study was Ti XVII. This was chosen because until recently, it was the most abundant impurity found in the DITE tokamak, occurring in concentrations of up to 1% of the electron density (UKAEA Culham Annual Report, 1980). Typical plasma conditions from the DITE tokamak were used in the calculations. The sixteenth ionization state of Titanium was the most highly populated state under typical DITE conditions of $n_e = 5 \times 10^{13} \text{ cm}^{-3}$ and $T_e = 300 \text{ eV}$ (from coronal calculations) and transitions of the

Ti XVII ion were investigated to see if there were any transitions of this species with wavelengths accessible to laser pumping.

Two forbidden transitions in the ground state of Ti XVII were chosen for the study (as they have suitable wavelengths accessible for laser pumping) after consideration of the data presented in Lawson and Peacock (1980), where complete listings of the $2s^k 2p^1 \rightarrow 2s^{k-1} 2p^{1+1}$ transitions in Li-like through to F-like ions of common impurity metals are contained. The $^3P_1 \rightarrow ^3P_2$ (3834.6 Å) and $^3P_0 \rightarrow ^3P_1$ (3370.6 Å) transitions were chosen. Details of these transitions are shown in Figure 7.1, and have transition probabilities of 220 s^{-1} and 450 s^{-1} respectively. Forbidden lines with transition probabilities as low as 10^2 to 10^4 s^{-1} have been observed in tokamak plasmas. Those lines with $\Delta n=0$ typically lie in the VUV to visible regions of the spectrum.

The following factors were considered in the feasibility study:

1. The power required to saturate the particular transition.
2. The population enhancement available.
3. The effect of the electron collision rates.
4. The number of photons detectable.
5. The laser characteristics required.

The collision strengths and transition probabilities calculated by Bhatia et al (1980) were used throughout, estimated by Bhatia to be correct to within a factor of two for the collision strengths and to within 20% for the transition probabilities.

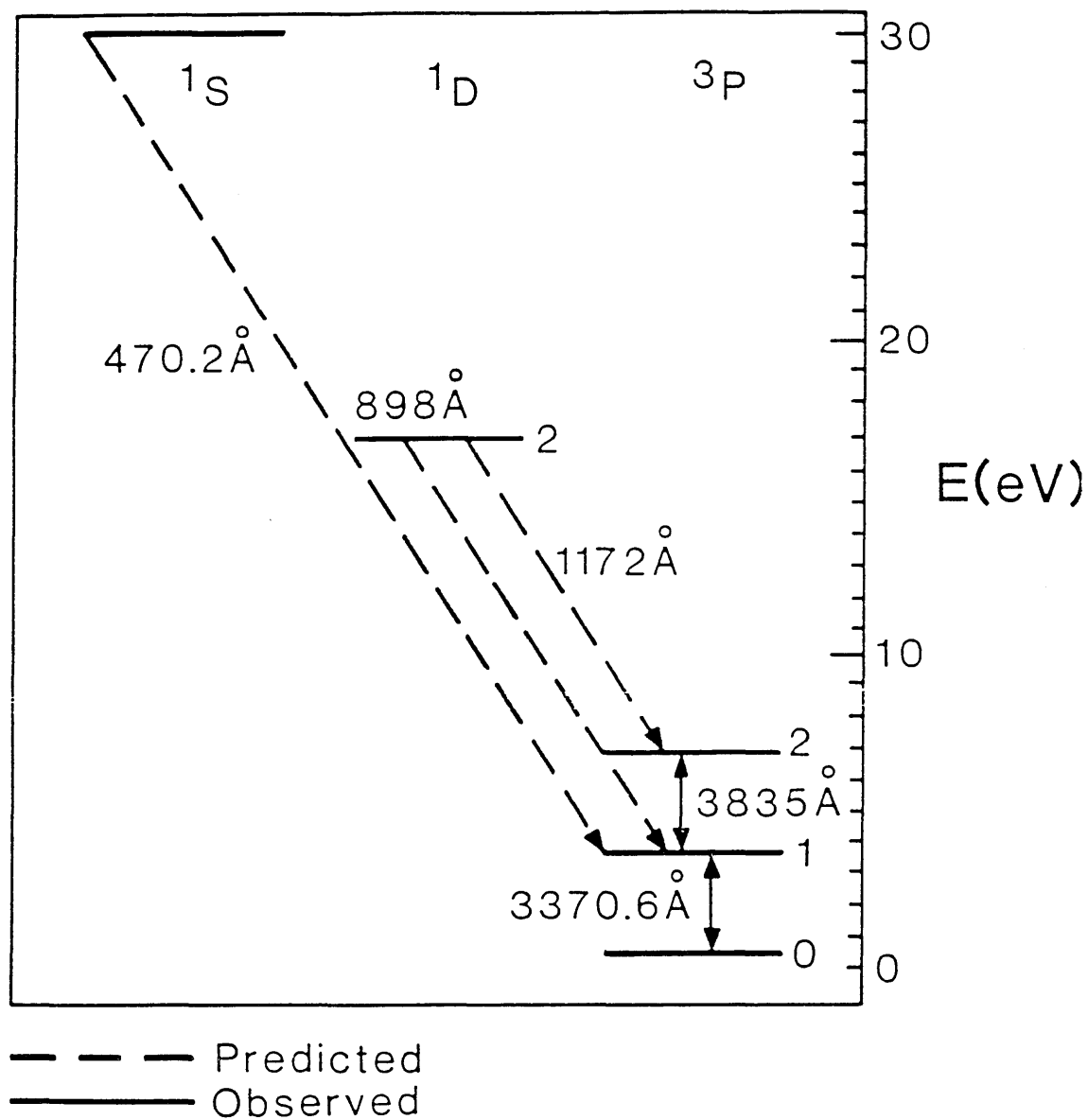


Figure 7.1 Transitions in the Ground State Configuration of Tl XVII.

7.2.2 Saturation Power Required. Consider a transition between two levels i and j , where i is the lower and j the upper level. In order to obtain substantial enhancement of the upper pumped level, the laser rate must exceed the total depopulating rate out of the upper level j .

Following the analysis of Burgess (1979) the power P (measured in Wcm^{-2} into the line width $\Delta\lambda$, into 4π solid angle) required to obtain a substantial enhancement of level j is:

$$P = \frac{8\pi hc}{\lambda^5} \frac{1}{A_{ji}(1-f_s)} \frac{f_s}{(K_{ji} + K_j + A_{ji} + (g_i/g_j)K_{ji})} \quad (7-2)$$

where g_i and g_j are the statistical weights of levels i and j respectively, K_{ij} and K_{ji} are the electron collision rates (s^{-1} per atom) from $i \rightarrow j$ and $j \rightarrow i$, A_{ji} is the spontaneous decay rate from the upper level and K_j is the total rate out of the upper level per upper level atom. All other symbols have their usual meanings. The saturation factor f_s is a measure of the degree of saturation, defined as:

$$f_s = \frac{n_j g_i}{n_i g_j} \quad (7-3)$$

f_s tends to unity as n_j/n_i tends to g_j/g_i , i.e. as the relative population of the two levels tend to the ratio of their statistical weights.

The electron collision rates were calculated by the author from the collision strengths tabulated by Bhatia et al (1980). Results from an alternative code constructed by Dr. H. Gordon of Strathclyde University were compared with the present results and found to agree to within 3%.

The Doppler widths (FWHM) of the transitions were calculated to be 0.65 Å for the $^3P_1 \rightarrow ^3P_2$ transition and 0.74 Å for the $^3P_0 \rightarrow ^3P_1$ transition at an electron

temperature of 300 eV (from Equation (6-1)). The Zeeman splitting for each transition was also calculated (from Equation (6-2) for a 2.7 T field). This was found to be ± 0.28 Å for the $^3P_1 \rightarrow ^3P_2$ transition and ± 0.22 Å for the $^3P_0 \rightarrow ^3P_1$ transition for the splitting of the sigma components (there being no splitting of the π components). Each separate Zeeman component will be Doppler broadened and to obtain the largest possible fluorescence signal, the laser linewidth must exceed the convolution of these two dominant broadening/splitting mechanisms.

Figure 7.2 shows the variation of the saturation power (calculated from Equation (7-2)) with electron density for a saturation factor f_s of 0.5 and into a linewidth 1.5 times the Doppler width to account for the relatively large Zeeman splitting. For typical tokamak densities (approx. $5 \times 10^{13} \text{ cm}^{-3}$) the power required to saturate the $^3P_1 \rightarrow ^3P_2$ transition is approximately 6 MWcm^{-2} and for the $^3P_0 \rightarrow ^3P_1$ is 5 MWcm^{-2} .

However, the main restriction on the saturation power is not the actual power required but the time taken to attain saturation. Equation (7-2) gives the power required to achieve saturation assuming that the laser pulse time duration is at least equal to the inverse of the spontaneous decay rate ($1/A_{j1}$).

For the work performed at the Balmer- α and - β wavelengths in hydrogen this assumption is correct, but not for forbidden transitions. In the case of the $^3P_1 \rightarrow ^3P_2$ transition, the laser would have to have a pulse length in excess of 4.6 ms, or 2.2 ms in the case of the $^3P_0 \rightarrow ^3P_1$ transition. Conversely, for a laser pulse of 10 μs the actual power required would increase by a factor of approximately 10^3 over the estimate given above.

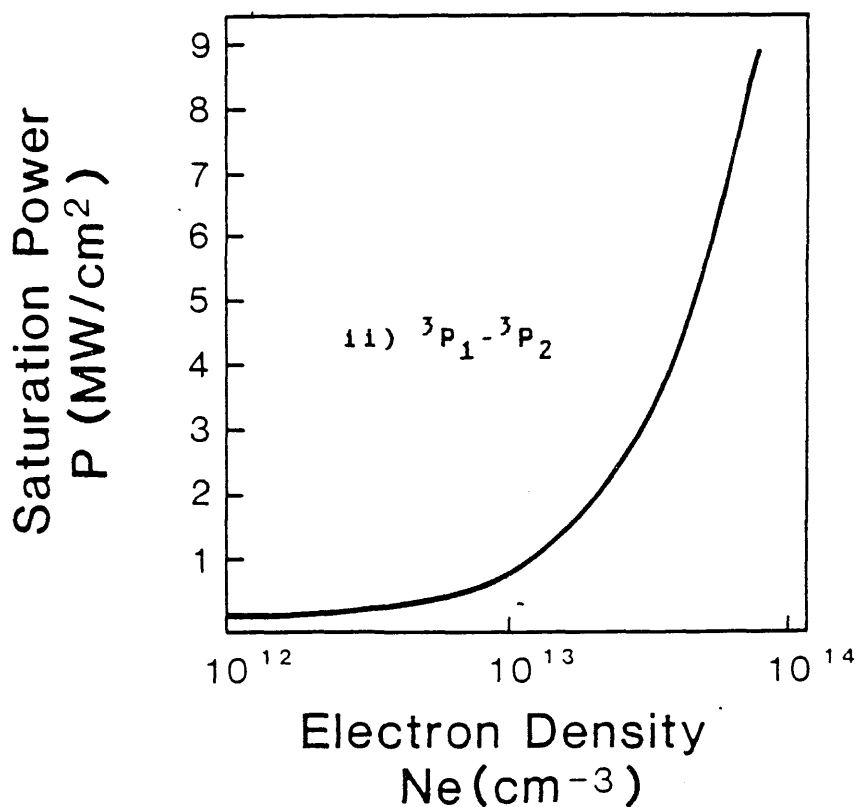
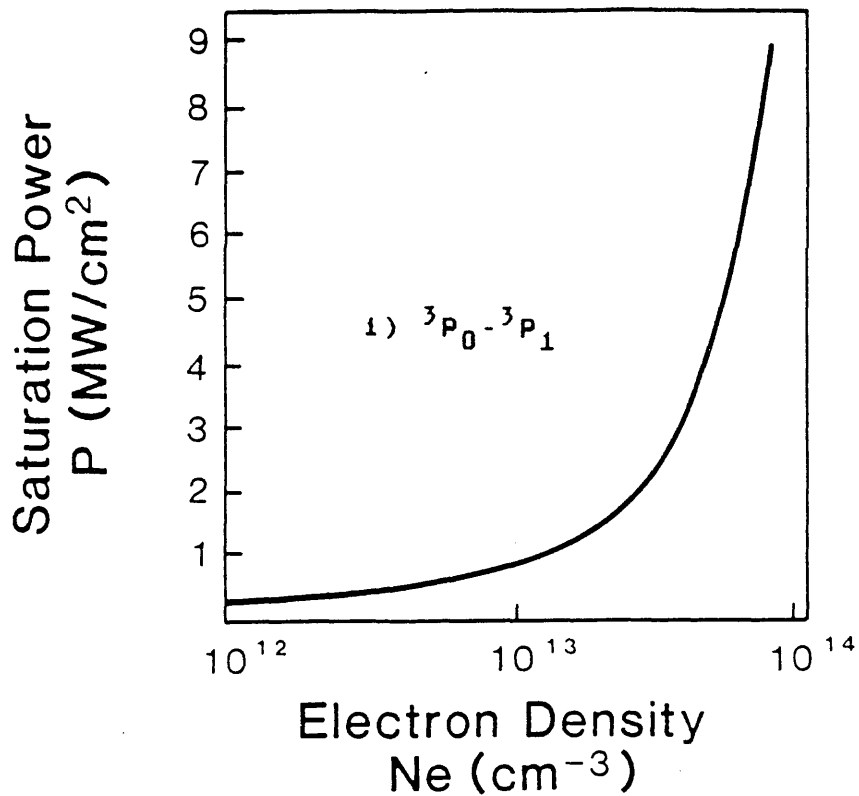


Figure 7.2 The Theoretical Saturation Power for the Tl XVII Transitions Under Investigation as a Function of the Electron Density.

This shows the power required to saturate the transitions in a time given by the inverse of their spontaneous decay rates at 300 eV into 1.5 time the Doppler width.

7.2.3 Population Enhancement. Using the level population densities calculated for a steady state from Bhatia et al (1980), the enhancement over background $\Delta n_U/n_0$ was calculated to be 0.15 for the $^3P_1 \rightarrow ^3P_2$ transition and 0.10 for the $^3P_0 \rightarrow ^3P_1$ transition for typical DITE operating conditions. Figure 7.3 shows this variation as a function of electron density for an electron temperature of 300 eV. This enhancement although small may be observable if the Ti XVII were present in a small region of the plasma only, in which case a very localised measurement may then be possible.

7.2.4 The Effect of Electron Collision Rates. The effect of the electron collision rates upon the enhancement expected must also be considered. In the case of the low temperature (<1 eV) z-pinch recombination plasmas discussed in Chapter 2, the collision rates out of the upper pumped level cause the initial enhancement to decay to a lower steady state value, where only the two pumped levels are out of equilibrium. However, in the present case, the electron collision rates out of the upper level (and indeed all the collisional and radiative rates associated with the pumped levels) are much smaller than for the previous case (by approximately 3 orders of magnitude). The sum of the non-laser rates out of the upper level K_U for both transitions studied is approximately $7 \times 10^4 \text{ s}^{-1}$ per upper state atom, compared to the saturating laser rate of $8 \times 10^7 \text{ s}^{-1}$ (for a power of 5 MWcm^{-2} into 1.5 times the Doppler linewidth). Therefore the initial enhancement will decay to e^{-1} of its initial value in a time given by $1/K_U = 14 \text{ ns}$. This time would essentially set the limit on the temporal resolution available with this technique. The laser pulse must also be powerful enough to equalise the populations well within this time for the enhancement to be observed.

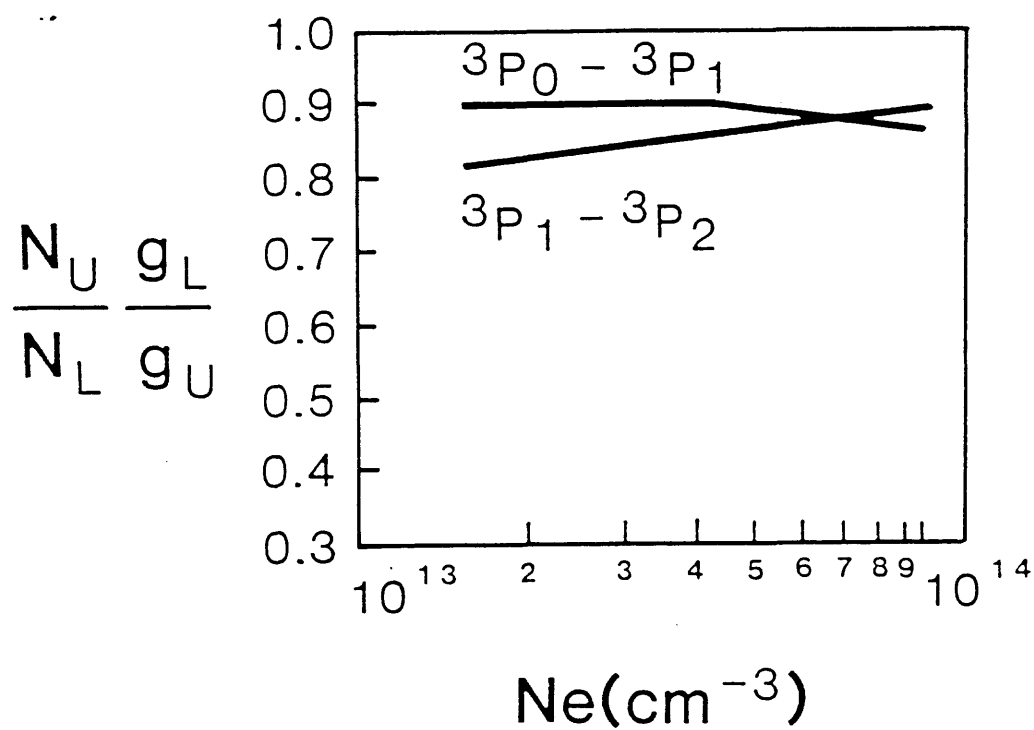


Figure 7.3 The Steady State Population Ratios as a Function of Electron Density.

The steady state populations were taken from Bhatia et al (1980) for $T_e=300$ eV.

7.2.5 The Number of Photons Detectable. The number of photons produced from a very fast saturating laser pulse (after which the fluorescence signal would decay exponentially in a time given by the collision rates out of the upper level) was calculated from the following expression from Gohil (1982):

$$N_{PH} = n A_{UL} V_F (\Omega/4\pi) \int_0^t (n_U(t) - n_U(0)) dt \quad (7-4)$$

where A_{UL} is the spontaneous decay rate, n is the efficiency of the detection system, V_F is the irradiated volume Ω is the solid angle subtended by the detection optics and $n_U(t) - n_U(0)$ is the change in the upper level density during the time t over which the enhancement is observed. If the fluorescence signal is observed over the time taken for the upper level to decay back to its pre-laser level ($1/K_U$) then the number of photons observed in this time Δt is:

$$N_{ph} = n A_{UL} V_F \Omega \Delta n \Delta t$$

where Δn is the change in the upper level population due to the saturating laser pulse.

To estimate the initial population density n_U^0 , the total titanium concentration was assumed to be $5 \times 10^{11} \text{ cm}^{-3}$, i.e. 1% of the electron density. From the abundances of the ion stages at 300 eV, the total density of the Ti XVII ion stage was calculated to be approximately $2 \times 10^{11} \text{ cm}^{-3}$. Assuming this population to be distributed only amongst the $^3P_2$, $^3P_1$ and $^3P_0$ levels, from the calculations of Bhatia et al (1980), the steady state population for the $^3P_2$ level was calculated to be $2 \times 10^9 \text{ cm}^{-3}$ and for the $^3P_1$ level, $3.2 \times 10^9 \text{ cm}^{-3}$. A detection efficiency of 20% and a collection solid angle of 10^{-2} steradians was assumed. K_U (the collision rate out of the upper level) for the $^3P_2$ level was calculated to be

$7.3 \times 10^4 \text{ s}^{-1}$ and for the $^3P_1$ level $7.8 \times 10^4 \text{ s}^{-1}$ per upper level atom. For an irradiated volume V_F of 2 cm^3 (giving good spatial resolution for a DITE minor radius of 26 cm) the $^3P_1$ - $^3P_2$ transition would yield 2.0×10^3 photons in a $1/K_U$ time of 14 μs , and the $^3P_0$ - $^3P_1$ transition would yield 1.8×10^3 photons in 13 μs .

7.2.6 Summary of Laser Characteristics. From the previous calculations, the laser required to saturate the two transitions of interest in Ti XVII would need to satisfy several stringent conditions. In order to saturate the chosen transitions in a time of 10 μs the output power must be approximately 5 GWcm^{-2} into a bandwidth of approximately 1.1 $\text{\AA}$ at a wavelength centred at 3370.6 or 3834.6 $\text{\AA}$. The time duration of the laser pulse would need to be as long as possible to maximize the number of detectable photons because of the small enhancement available. If a laser with pulse length longer than the inverse of the spontaneous rate out of the upper level were available (milliseconds) then the required power drops to 6 MWcm^{-2} .

As an example of currently available laser powers, the maximum output power reported for a commercial tripled Nd:YAG system is 50 MWcm^{-2} at $\lambda=3550 \text{ \AA}$ (Laser Focus/Electro-Optics Buyers' Guide, 1985). However, pulsed dye lasers would give the desired tunability.

The flashlamp pumped dye laser described in Chapters 3 and 4, and used in the experimental fluorescence scattering work described in this thesis may be a suitable candidate for this type of work, with its relatively long pulse length, tunability and high power output. The maximum power output obtained for the Coumarin dye ($\lambda=4861 \text{ \AA}$) was 1.1 MWcm^{-2} into a 5 $\text{\AA}$ bandwidth and a time duration of 400 ns. The usual laser configuration could be modified, i.e. by the addition of

more amplification stages and an increase in the xenon pressure and firing voltage. Focussing the laser beam down would give an increased power density too, although this may adversely affect the stray light level.

A suitable dye for pumping the $^3P_1 \rightarrow ^3P_2$ transition ($\lambda=3834.6 \text{ \AA}$) would be p-Terphenyl dissolved in DMF, which lases from 3700-4100 $\text{\AA}$ when pumped by a flashlamp source (Furumoto and Ceccon, 1970). A maximum power output of 150 KW is quoted for this system (Laser Focus/Electro-Optics Buyers' Guide, 1985) for a pulse length of 200 ns. The maximum power output for a tunable dye system which would lase in this region is quoted to be 20 MW for a commercial excimer or YAG pumped system (Laser Focus/Electro-Optics Buyers' Guide, 1985). For lasing action at $\lambda=3370.6 \text{ \AA}$, the wavelength of the $^3P_0 \rightarrow ^3P_1$ transition, BBQ dye dissolved in DMF lases in the range 3350-3550 $\text{\AA}$, again when pumped by a flashlamp (Exciton Dye Catalogue, 1985). However these laser intensities are simply not large enough to produce saturation of the transitions discussed here.

7.3 Conclusions.

Two forbidden transitions of Ti XVII (the $^3P_1 \rightarrow ^3P_2$ and $^3P_0 \rightarrow ^3P_1$) were studied to see if it would be practical to perform a laser induced fluorescence scattering experiment to measure directly the density of such a highly ionized species. The wavelengths of the two transitions were 3834.6 and 3370.6 $\text{\AA}$ respectively. The main problems highlighted by the feasibility study were the low enhancement over background expected (10-15% over background) and the high laser powers required to saturate the transitions (approximately 5 GWcm^{-2} in a pulse of 14 us, falling to 6 MWcm^{-2} sustained over a few milliseconds).

The low enhancement in itself may not be a problem if the Ti XVII is localised in a small region, i.e. the plasma core, and because the long decay time of the fluorescence allows the signal to be integrated. The small number of photons predicted (approximately $1.6-2.0 \cdot 10^3$) could be increased by increasing the irradiated volume or the detector efficiency for example. The major obstacle to the use of fluorescence scattering techniques for impurity ion diagnostics is the extremely high laser intensities required.

Chapter Eight

SUMMARY, CONCLUSIONS AND SUGGESTIONS FOR FURTHER WORK.

8.1 Summary and Conclusions.

The work described in this thesis has a common link, although involving two different sets of plasma conditions. This link is the use of the technique of laser induced fluorescence scattering.

The two different regimes investigated were:

i) the recombination phase of a low temperature ($0.19 < T_e < 1.03$ eV) z-pinch device in hydrogen. Under these conditions the ground state is effectively decoupled from the excited levels. This work involved understanding the discrepancies between the experimentally observed and the theoretically predicted H- α and collisionally induced H- β fluorescence signals.

ii) high temperature ($T_e > 100$ eV) fusion-type plasmas where the ground state was coupled (via the electron collisional rates) to the excited states. This work mainly concerned the use of the fluorescence scattering technique as a diagnostic for the determination of the spatially resolved neutral and impurity density. The neutral density measurement was carried out on the reverse field pinch (HBTX1A) at the UKAEA Culham Laboratory and a feasibility study on a highly ionized species was carried out assuming typical operating conditions for the DITE tokamak, also situated at Culham Laboratory.

8.1.1 Low Temperature Work. Previous work in the recombination phase of one particular, well diagnosed z-pinch device, had shown that considerable discrepancies between the experimental data (as observed by Burgess et al, 1980) and the theoretical predictions existed.

In an attempt to account for these discrepancies the following new features were incorporated into the CR code:

i) a new set of electron collision rates (Vriens and Smeets, 1980) which were typically a factor of 3 smaller than the previously accepted set (Johnson, 1972).

ii) extra processes involving LTE and steady state densities of H^- and H_2^+ ions which could populate the $n=2,3$ and 4 levels. The presence of other species in the plasma such as H_2 and H_3^+ and their effect upon the observed discrepancies was discounted on the experimental evidence of Gohil (1982).

This, plus the addition of error bars on the theoretical predictions (to reflect the uncertainty in the diagnosed electron temperature) had the effect of eliminating or reducing some of the discrepancies previously observed.

However, two anomalies still could not be explained by this:

i) the collisionally induced $n=4$ plateau enhancement is a factor of 1.5 times larger than the theory, even when the original rates of Johnson are reduced by a factor of 10. The $n=3$ plateau enhancement at the higher densities ($n_e > 2 \cdot 10^{14} \text{ cm}^{-3}$) is also significantly larger than theory.

ii) the absolute $n=2$ density measured by Burgess et al (1980) which is larger than that predicted by theory especially for the higher densities (for $n_e > 8 \cdot 10^{14} \text{ cm}^{-3}$ the data is an order of magnitude higher). Addition of the extra rates into the $n=2,3$ and 4 levels increased the $n=2$ density by less than 10%.

Two experiments were carried out on this plasma source to

investigate the role of H^- ions in keeping the $n=2$ density anomalously high. The first experiment involved using a high power $H-\alpha$ laser to deplete the H^- ion population (via photodetachment) and to observe the temporal behaviour of the characteristic fluorescence plateau. The second involved removing the H^- ions from the plasma with a high power ruby laser pulse (which would also photoionize the $n \geq 3$ levels) and observing the behaviour of the $H-\alpha$ and $H-\beta$ emission as it returned to its initial background level.

From the first set of experiments, an anomalously high $H-\alpha$ plateau was clearly visible, confirming previous work on the same plasma device. At delay times < 70 us, the plateau enhancement was observed to decay whilst the $H-\alpha$ laser was still saturating the transition. It was concluded that the $n=2$ level was populated by a process involving a species which is removed by the action of the high power $H-\alpha$ laser beam. The most likely candidate is H^- , populating $n=2$ via mutual neutralization. If this process is responsible, then to be consistent with the work presented in Chapter 2 either the H^- density is greater than predicted assuming LTE, or the cross section into the $n=2$ level must be equal to or greater than that into the $n=3$ level.

In the second set of experiments, the $n=3$ level was depleted by photoionization by a ruby laser beam and was observed to repopulate in a time which was slower than that given by the electron collision rates out of the $n=3$ level. The $n=3$ level repopulated in a time faster than or approximately equal to the H^- recombination time at the earlier delay times, but slower than this recombination time at the 50 and 60 us delay time data. The $n=4$ level was observed to repopulate in a time faster than the H^- recombination. These results suggest that the $n=4$ and $n=3$ levels (for the 50 and 60 us delay times) are not significantly populated by processes involving H^- ions. For the $n=3$ results concerning the 20, 30 and 40 us delay

times, either this level is partially populated by processes involving H^- ions, or the collision rates out of this level are factors of approximately 2 to 6 lower than given by Vriens and Smeets. This latter conclusion would be consistent with earlier work.

A series of $H-\alpha$ fluorescence measurements were also made in the recombination phase of another z-pinch device, which was in general slightly hotter and denser than the original z-pinch. Similar discrepancies were found as regards the $H-\alpha$ peak and plateau enhancements. The anomalies were thus shown not to be a characteristic of a particular device, although the measured peak decay rate agreed with that predicted using the rates of Vriens and Smeets.

8.1.2 High Temperature Work. A spatially resolved neutral density profile was obtained using, both $D-\alpha$ and $D-\beta$ fluorescence scattering in the reverse field pinch HBTX1A. The profile was slightly asymmetric, the neutral density being higher towards the outer wall.

Addition of a vertical field caused the observed neutral density to decrease, this effect being more pronounced for the most central position in the plasma where the density was observed to decrease from $7.28 \pm 1.79 \times 10^{11}$ to $3.37 \pm 0.74 \times 10^{11} \text{ cm}^{-3}$ for the 4 cm resolution data points. This result confirmed previous work (Radcliffe, 1984) that the addition of such a field caused the usually unsymmetric plasma to shift away from the outer wall to a more central position, and also showed that the origin of the neutrals in the plasma was primarily from plasma-wall interactions such as the dissociation of the D_2^+ ions originating from the ionization of D_2 from the wall.

Neutral density profiles were generated from a transport code (the AURORA code) and compared to the experimentally determined profile. The only good match was obtained for the region between the plasma centre and the inner wall,

for an input energy of the edge neutrals of 5.5 eV. It was postulated that the inclusion of charge exchange with impurities in the AURORA code would decrease the disagreement between the data and the output from the code for the region near to the outer wall.

The fluorescence enhancements were also measured and compared with the predicted values from the CR code. These agreed to within experimental error.

Finally, an investigation into the use of the fluorescence scattering technique on highly ionized species as a diagnostic for the plasma core was carried out on two particular transitions of Ti XVII ($^3P_1 \rightarrow ^3P_2$ at 3834.6 Å and $^3P_0 \rightarrow ^3P_1$ at 3370.6 Å). It was concluded that favourable plasma conditions and careful design of the experimental system would be required because of the low enhancement expected (10-15% over background) but that the major problem was due to the long laser pulse and high power output required from the laser. For a pulse duration of 14 ns a power of 5 GW cm⁻² into a bandwidth of 1.1 Å is required for the species studied.

8.2 Suggestions for Further Work.

8.2.1 Low Temperature Work. It would be interesting to repeat and extend the n=2 density measurements to check the anomalously high population obtained over a wider range of plasma conditions (previous work only covered the 30-50 ns delay times).

To check the hypothesis that H⁻ ions are indeed responsible for the high n=2 density, it would also be interesting to monitor the n=2 population (in absorption) whilst at the same time depleting the H⁻ density with a ruby or Nd:YAG laser.

Much work involving the higher lying levels in hydrogen could be performed with sensitive detection equipment. A

7.5 cm aperture Fabry-Perot has been acquired by the Spectroscopy Group at Imperial College. This has more than 30 times the light throughput available than the standard Hilger Monospek monochromator. In particular, the $n=3,4$ and 5 risetimes whilst saturating the H- α transition could be remeasured using this system with a short cable and a fast oscilloscope. Very little work has also been done on the H-He (5% H by volume) plasma because of the low hydrogen line intensities involved, and this work could also be extended.

8.2.2 High Temperature Work. The fluorescence scattering technique may be extended as a diagnostic for fusion-type plasmas.

As described in Section 6.4.3, if the laser is extinguished whilst still saturating the transition, the fluorescence signal will decay in a time given by the collisional rate out of the upper level into the lower level, and if the collision coefficient is known then the local electron density may be calculated from the measured decay time. For the range of densities in HBTX1A ($1\text{--}3 \times 10^{13} \text{ cm}^{-3}$) the D- β transition gives the largest decay time variation across the range of electron densities plus decay times that are not too small to measure ($2\text{--}10 \text{ ns}$). This experiment was attempted after the neutral density profile had been obtained, but was unsuccessful for reasons explained in Section 6.4.3.

Another application of this technique would be in the measurement of the local ion temperature. The laser could be line narrowed (addition of an extra Fabry-Perot etalon in the oscillator cavity would achieve this) and scanned across the fluorescence linewidth. If the fluorescence line was mainly Doppler broadened a spatially resolved temperature could be obtained from measuring the width of the line. This would be independent of the electron density profile but relies on the reproducibility of the plasma. A better way of achieving the same result would

be to ensure that the laser bandwidth was sufficiently large to excite all the available neutrals in the irradiated volume (this may not be easy if there is a large Zeeman component) and observe the entire fluorescence line for one shot by using an Optical Multichannel Analyser (OMA).

With a carefully designed experimental arrangement, it may be possible to obtain with one plasma shot, not only the local neutral density measurement but also the local electron density and ion temperature.

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