RADIATION TRANSPORT AND SHORT PULSE INTERACTION IN LASER IRRADIATED TARGETS

A thesis submitted on application for admission to the degree of Doctor of Philosophy at the University of London by Michael John Edwards, Imperial College of Science, Technology and Medicine, October 1989.
ABSTRACT

Experimental measurements of radiation transport through plastic foils up to 10 μm thick obtained using time resolved XUV spectroscopy in the 10 - 70 Å spectral wavelength region are presented. An intense radiation source was produced by overcoating one side of the target with 0.1 μm of gold and irradiating it with green laser light at an intensity close to $10^{14}$ Wcm$^{-2}$. Detailed modelling of the results obtained with a 1D radiation hydrodynamics code show good overall agreement with the experiment. It is shown that radiation transport dominates the energy transfer in the targets and that the atomic level populations are maintained closer to LTE when an ambient radiation field exists. The effect of edge shifting, edge smearing, outer shells opacity and the radiation pulse's temporal and spectral shape on the radiation transport are investigated.

The production of hot (400 eV) high electron density ($\approx 10^{23}$ cm$^{-3}$) aluminium plasmas by the irradiation of solid targets with high intensity ($\approx 10^{17}$ Wcm$^{-2}$) 3.5 ps, KrF laser pulses is briefly described. The experimental conditions are simulated using a 1D Lagrangian hydrocode. Good agreement is obtained between experiment and the modelling. Preliminary predictions are made that hot ($> 0.5$ keV) plasma might be produced at solid density by a 300 fs, 1 μm laser pulse.
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To Deborah, my family and other animals.
CHAPTER 1

1.1 INTRODUCTION

When a high power laser beam is focussed onto a solid target, a dense plasma is formed in which, predominantly, soft X-ray or XUV radiation (100 eV to 1 keV) is produced. These types of plasmas are very important for research in Inertial Confinement Fusion (ICF), dense plasmas, X-ray Lasers (XRL) and even stellar atmospheres.

The X-rays are produced in the form of lines or bound-bound (b-b) transitions superimposed on a continuous background of free-bound (f-b) (recombination radiation) and free-free radiation (bremsstrahlung). The relative importance of the contributions depends mainly on the atomic number (Z) and degree of ionization of the target material and the plasma temperature and density. The fraction of the laser energy converted into X-rays increases as Z increases. This is because the electrons experience larger accelerations due to the presence of larger nuclear charges in the plasma. Consequently, the rate at which the electrons radiate is increased.

Most of the radiation is produced in the region just above the critical plasma density (the maximum electron density to which the laser light may propagate). This is because, although the density rises rapidly past this point, the temperature and thus degree of ionization falls sharply and the rate of production of X-rays (generally) increases with the square of the plasma density. Consequently when shorter wavelength laser light is used the laser to X-ray conversion efficiency is increased because the laser light can propagate to a higher electron density.
However, short wavelength lasers are advantageous for major research programmes such as ICF and XRL because the laser energy may propagate to a higher electron density where it is far more efficiently absorbed. Recent measurements have shown an X-ray conversion efficiency of 80% on gold targets when irradiated by 0.26 μm laser light (Kauffman 1986, Kodama et al 1986). In addition, instabilities generating hot electrons are inhibited by shorter wavelength laser light.

Under these conditions, because a significant fraction of the absorbed laser energy is converted to X-rays, radiation transport becomes a major energy transport mechanism. In the past, radiation transport has been largely ignored in laser plasmas. This is mainly because the fraction of the absorbed laser energy converted into soft X-rays in laser target interaction experiments has generally been of the order of or less than a few percent. This results from the fact that previous experiments have been chiefly concerned with laser targets made of low Z materials and have used long laser wavelengths (> 0.5 μm).

This thesis describes both experimental measurements of radiation transport into the cooler denser regions of laser targets using time resolved XUV spectroscopy for the first time and an in depth study of the radiation transport using a radiation transport code coupled to a 1D Lagrangian hydrodynamics code. The results of the simulations and experiment are compared.

In addition, the modelling of ultra short (3 ps), high intensity (> 10^{16} \text{ Wcm}^{-2}), KrF laser pulses with solid targets using a 1D Lagrangian hydrocode is described. The case when a low intensity
amplified spontaneous emission (ASE) prepulse forms a large, cool, underdense plasma with which the high intensity, short pulse interacts is also considered. This modelling was performed in support of an experiment in which the author was involved and which is briefly reviewed. Finally, on the basis of the results of this modelling, preliminary predictions of plasma conditions that may be expected to be produced by the Imperial College (IC) short pulse laser system, currently under construction, are presented.

1.2 PREVIOUS WORK: Experimental Studies Related to Radiation Transport.

Experimentally, very little has been done to study radiation transport into the cooler, denser regions of laser targets. There are two reasons for this. Firstly, radiation transport has not generally been considered important because of low Z target materials and long laser wavelengths. Secondly, most of the radiative energy is produced in the soft X-ray or XUV energy region and suitable diagnostics to analyse radiation of these wavelengths have only recently become available. Consequently, the majority of experiments have been concerned with radiation transported out of (emitted from) plasmas just above critical density where most of the radiation is produced. Time resolved X-ray crystal spectroscopy below 10 Å (above 1 keV) and, more recently, with the development of flat field spectrometers, time resolved VUV spectroscopy above 100 Å (below 150 eV) have been used.

Few experiments have been reported in the literature which have
directly attempted to measure radiation transport in the soft X-ray region. Mizui et al (1981) studied radiation emitted from the rear of laser irradiated gold and aluminium foils using time integrated XUV spectroscopy. They concluded that radiation conduction played an important role in the energy transport in high Z plasmas. A similar conclusion was reached by Mochizuki et al (1987) who analysed the X-rays from a gold source laser plasma transmitted through very thin aluminium and gold foils using diodes and a time resolved transmission grating spectrometer. A shift of the aluminium L-edge due to ionization was also inferred.

Contradictory results to those of Mizui et al (1981) were found by Nishimura et al (1983). They found that mass ablation rate data inferred from diode measurements on gold targets agreed best with electron thermal conduction models. In contrast, a simple model based on radiative ablation gave a mass ablation rate at least an order of magnitude higher, in agreement with the mass ablation rate inferred by Mizui et al (1981).

There are several other categories of experiments which do not directly measure radiation transport but which are related to this topic. Firstly, X-ray preheat experiments assess the extent of the heating of the rear side of foil targets, before a hydrodynamic shock wave breaks through, due to hard X-ray radiation depositing energy at the back of the foil. The technique most extensively used is the optical imaging of the rear side of, in particular, laser irradiated aluminium foils at two wavelengths (McLean et al 1980). This technique was extended by Ng et al (1985) who used a prism to obtain a continuous optical spectrum from the rear of laser heated foils. More recent preheat
measurements have been made using novel multilayer Schwarzschild microscopes to obtain time resolved images of the rear side of laser irradiated gold and aluminium foils in the UV at 10, 40 and 80 eV (Benattar 1987).

Secondly, X-ray absorption spectroscopy experiments above 1 keV, which are distinct from transport experiments because the re-emission of photons is almost certainly not important, have been used to analyse laser driven implosion targets (Yaakobi et al 1979, Hauer et al 1986) and laser heated foils (Chenais-Popovics et al 1987). This technique has also been used recently to assess the opacity close to local thermodynamic equilibrium (LTE) of tamped aluminium foils, heated directly by X-rays, in the region of the aluminium K-edge (Davidson et al 1988). Other experiments have studied ionic structure (Hall et al 1988) and the effects of free electron degeneracy and continuum lowering (Bradley et al 1987, Riley et al 1989) in strongly coupled, partially degenerate, aluminium and chlorine plasmas respectively.

Finally, radiation confinement has been studied in laser heated cavities for ICF applications (Sigel 1987, 1988).

1.3 THEORETICAL STUDIES OF RADIATION TRANSPORT IN LASER PRODUCED PLASMAS.

Radiation transport modelling has been applied to laser plasmas on numerous occasions and is reviewed by Lee (1982). The majority of the modelling has been concerned with the prediction and interpretation of X-ray emission spectra originating from around the
critical surface (Coulambant et al 1975, Duston and Davis 1980, 1981, Lee 1981). This is mainly because the majority of experimental efforts were also concentrated on emission spectroscopy. However, radiation transport as an energy transport mechanism has been included in laser hydrocodes in some cases.

X-ray conversion efficiencies and global spectra predicted by the 2D Lagrangian hydrocode, LASNEX, including multigroup diffusion have been compared with experimental measurements for various target materials (Haas et al 1977, Shay et al 1978, Rosen et al 1979, Mead et al 1983, 1984).

A multigroup diffusion radiation hydrodynamics code (Ramis, Schmalz and Meyer-ter-Vehn 1986) as well as self similar techniques (Pakula and Sigel 1986, Sigel 1987, 1988, Tsakiris et al 1988) have been used to investigate radiation confinement in laser heated cavities.


Distinct from the radiation diffusion approximation, Rose and Evans (1983) have used a forwards/reverse approximation to the equation of radiative transfer in a 1D Lagrangian hydrodynamics code to study the effects of radiation transport on plasma gradients in laser driven microballoon implosions.

The effects of radiation on plasma gradients, spectra and preheat in laser irradiated aluminium and carbon foil targets has been investigated in detail by Duston et al (1983) and Duston, Clark and Davis (1985)
using a probability of escape technique (Aruzese et al 1980) for the 
radiation transport in a 1D hydrocode.

Finally, multigroup radiation transport in a 1D Lagrangian hydrocode 
has been reported in the investigation of X-ray conversion efficiencies 
as a function of target material (Mochizuki et al 1986).

1.4 AUTHOR'S CONTRIBUTIONS TO THE STUDY OF 
RADIATION TRANSPORT IN LASER PLASMAS.

The author has designed and performed an experiment to measure 
radiation transport into the cooler, denser regions of plastic foil targets 
of various thicknesses subjected to an intense, pseudo Planckian 
radiation source. The radiation transported through the plastic foils was 
analysed using time resolved XUV spectroscopy in the 10 to 70 Å 
spectral wavelength region. These are the first direct time resolved 
XUV measurements of radiation transport which utilise the very high 
spectral resolution properties of a flat field XUV spectrometer 
(250 mA). Higher order contributions were reduced to a negligible 
level in the experiment through the addition of a grazing incidence 
X-ray mirror to the instrument.

In contrast, previous measurements have used XUV diodes (Nishimura 
et al 1983, Mochizuki et al 1987) whose spectral resolution is 
$\sim$ 100 eV. The time resolved measurements (Mochizuki et al 1987) 
have used transmission grating spectrometers which suffer from 
extremely efficient multiple order contributions and have a comparatively
poor spectral resolution ~ 4 Å.

Low Z targets have been used by the author because the atomic physics and radiation transport in higher Z plasma is extremely complex. Therefore, experiments using high Z plasmas seem unlikely to shed light on the underlying physical processes or enhance the understanding of radiation transport, particularly when very contradictory experimental results are obtained (between Mizui et al 1981 and Nishimura et al 1983).

The experimental results have been analysed in detail by the author with a multigroup radiation transport model coupled to a 1D Lagrangian hydrodynamics code. The radiation transport algorithm was written by Rose and Evans (1983). The radiative opacity is calculated using an average atom (AA) model assuming LTE and has been developed by the author. Unlike the majority of previous work, group averaged Planck mean opacities are calculated in-line with the hydrodynamics for 116 photon energy groups at every time step in every fluid cell.

The radiation hydrodynamics code has been used by the author to obtain an in depth understanding and fully quantitative description of radiation transport in the radiatively heated foils. This work is distinct from previous reported investigations (Rose and Evans 1983, Duston et al 1983, Duston, Clark and Davis 1985) because it considers a low Z target subject to a high intensity, pseudo-Planckian radiation field. Also, most of the radiation in the plastic travels into the cooler, denser regions towards the rear of the sample instead of out into the vacuum as in conventional laser plasma interaction calculations.
Finally, non-LTE effects on the atomic level population at various positions in the foils are quantitatively assessed by the author using a zero dimensional non-LTE AA model written by Dr. S. J. Rose. Particular emphasis has been placed on the effects of the ambient radiation field in the plasma. To the best of the author's knowledge, no details exist in the literature on this topic for targets subject to these conditions.

The thesis is presented as follows.

In chapter 2, elementary radiative transfer is discussed. All approximations and assumptions are verified when possible.

The radiation transport experiment, diagnostics and experimental data is described in chapter 3.

Chapter 4 gives a detailed step by step account of the radiation transport model and development of the opacity code. The application of the model to the experiment is discussed in detail.

A detailed account of the radiation transport simulations is presented in chapter 5 and the results are compared to the experimental data. Material heating, outer shells opacity scaling, radiation pulse shape (spectral and temporal) and non-LTE effects are quantitatively discussed.

Chapter 6 describes an experiment carried out using high intensity 3.5 ps KrF laser pulses. Detailed hydrocode modelling is presented for cases with and without prepulses and preliminary predictions for the
300 fs Imperial College laser are discussed.

The results and conclusions are summarised in chapter 7.

All units are SI unless otherwise stated.

Angstroms are denoted thus, A.
CHAPTER 2

ELEMENTARY THEORY OF RADIATIVE TRANSFER

It is not the purpose of this section to give a full account of the theory of radiative transfer which can be found in numerous, excellent texts (see for example Mihalas and Mihalas 1984). Here, beginning with the equation of radiative transfer, without any specific details of the material opacity, the equations forming the basis of the radiation transport model are developed. Then the individual sources which contribute to the material opacity are described and those which are not relevant are identified and eliminated. First, however, several simplifying assumptions are made in addition to those of unpolarized, incoherent radiation.

2.1 SIMPLIFYING ASSUMPTIONS

2.1.1 Local Thermodynamic Equilibrium

Complete thermodynamic equilibrium (TE) holds in an isolated system when the material and radiation equilibrate to a common temperature. In TE each process is in detailed balance with its inverse so that a steady state obtains. Furthermore, the radiation field is then isotropic and Planckian everywhere. In addition, the distribution of particle velocities and of particles amongst energy states is given by Maxwell-Bolzmann statistics. Hence, although radiative processes are not negligible, the particle distributions are the same as if only collisional processes were occurring. If a system exists, in which the temperature may vary but at a given point the particle distributions are given by Maxwell-Bolzmann statistics at the local kinetic temperature,
then local thermodynamic equilibrium (LTE) is said to exist at that local temperature. This definition of LTE makes no prerequisites on the nature of the radiation field. However, beginning from the Boltzmann expression for the distribution of particles amongst energy states, a modified form of Kirchhoff's law may be obtained (Armstrong and Nicholls 1972; Kirchhoff's law is sometimes stated as the condition for LTE). Hence, if LTE exists the radiation emission coefficient can be obtained directly from the modified Kirchhoff law, which will be described later.

From the definition of LTE, it is clear that if collisional processes are dominant in a system then LTE will obtain. However, this is not a prerequisite for LTE. Provided the radiative transitions also occur at their equilibrium values (e.g. Planckian radiation field) then the radiation will tend to drive the material towards (or help maintain it in) a state of LTE. If the collisional processes are not dominant and the radiative processes are not in detailed balance then the radiation will tend to force the system away from LTE. An example of this is the low density, transparent coronal plasma which is formed in laser target interaction experiments. Collisional processes occur more infrequently than radiative recombination or spontaneous emission and because the plasma is very transparent (radiative ionization and excitation are inefficient) the ionic ground states become more populated than in LTE.

In the case of the comparatively cool, dense, supercritical plasma which is considered in the modelling and experiment described in chapters 3, 4 and 5, the assumption of LTE is made on the basis that collisional processes are likely to be dominant. Although no quantitative
justification is given here, a further discussion taking into account the heating radiation is presented in chapter 5 where it is shown quantitatively that LTE is a good approximation in the dense foil plasmas.

2.1.2 Unit Refractive Index

The dispersive nature of a plasma can add considerably to the radiation transport problem. Modifications must be made both to the transfer equation and to the material opacity (Cox 1968). However, because the vast majority of radiation produced in a laser plasma is in the soft X-ray region, approximating the refractive index as unity is valid to a high degree of accuracy.

The refractive index for a plasma, in the absence of any magnetic fields, due to free electrons, is given by (Hughes 1979)

\[
\eta = \left( 1 - \frac{\omega_e^2}{\omega_\nu^2} \right)^{1/2}
\]

2.1.1

where \(\omega_e\) is the angular electron plasma frequency proportional to the square root of the electron density and \(\omega_\nu\) the radiation angular frequency. If \(\omega_e/\omega_\nu << 1\), equation 2.1.1 may be rearranged to give

\[
1-\eta = 5 \times 10^{-4} \frac{Z^* \rho}{A} \frac{\omega_\nu^{-2}}{
\]

where \(\omega_\nu\) is in keV, \(Z^*\) is the average degree of ionization, \(A\) the relative atomic mass, \(\rho\) the density in g cm\(^{-3}\), \(Z^* \rho/A\) is the electron density in mols cm\(^{-3}\) and use has been made of the electron density dependence of \(\omega_e\). Even for fully ionized plasma at solid density (1 g cm\(^{-3}\)) the above expression shows \(\eta\) to be different from unity by less than 2.5% for 100 eV photons. Therefore, because most of the radiation is more energetic and the above conditions are extreme for
the radiatively heated foil plasmas, the approximation of unit refractive index will not result in any significant error in the material heating.

The case of refraction due to bound electrons is not considered here. A detailed discussion can be found elsewhere (Crowley 1985).

2.1.3 Scattering

By far the most significant scattering process in a laser plasma is Compton scattering. In the classical limit, valid in the radiatively heated foil plasmas, this is equivalent to Thomson scattering in which the photon energy is essentially left unchanged. The Thomson cross-section has the value of $6.65 \times 10^{-25}$ cm$^2$ per electron and is frequency independent. At solid density ($1 \text{ gcm}^{-3}$) for fully ionized plasma, this corresponds to an opacity of order $10^{-1}$ cm$^2$g$^{-1}$. The next smallest contribution to the opacity in the foil plasmas is in general due to f-f transitions. For the same conditions as above for photon energies $\ll 1 \text{ keV}$ and plasma temperatures $\ll 1 \text{ keV}$, the f-f opacity is $\gg 10^1$ cm$^2$g$^{-1}$, two orders of magnitude larger than the scattering opacity. Therefore, scattering can be ignored.

2.1.4 Radiation Energy Density

It is shown here that the contributions to the total energy and momentum of the system due to the radiation can be neglected. In this way, the radiation is considered as a mechanism by which energy is redistributed within the system and can be treated as such in the hydrodynamics equations. In this section, this simply reduces the number of radiative quantities that must be considered. However, it considerably reduces the complexity of the modelling which will be described in chapter 4.
It is only necessary to consider the internal energy of the plasma and the radiation field. This is because the radiation and material pressures are of the same order as the respective energy densities. The material energy density is given by $U_m = C_v T$ and the radiative energy density is given by $U_r = 4\sigma T^4/c$ where $C_v$ is the specific heat per unit volume and $\sigma$ is Stefan's constant. For $T > 1$ these quantities converge to equality before diverging again. The temperature at which equality occurs is given by

$$T_{\text{equ}} = \left(\frac{4C_v c}{\sigma}\right)^{1/3} \approx 0.5 \text{ keV}$$

Therefore, because the temperature is always less than 0.5 keV in the foil plasmas, the largest value of $T$ will give the largest radiation to matter energy density ratio. Taking the typical foil values (4.7.2) $T \approx 100 \text{ eV}$, $\rho \approx 0.1 \text{ gcm}^{-3}$, $U_r/U_m \ll 5 \times 10^{-3}$. Hence, ignoring the contribution of the radiative energy and pressure in the fluid dynamics equations is an extremely good approximation. It has been assumed above that the material and radiation temperatures are equal. Even if this were not so, the radiation temperature would have to be $\sim 200 \text{ eV}$ to increase the above ratio to $5 \times 10^{-2}$.

## 2.2 EQUATION OF RADIATIVE TRANSFER

In general form, in an inertial frame, the equation of radiative transfer may be written as follows

$$\rho^{-1} \left[ \frac{\partial}{\partial t} + \frac{\partial}{\partial s} \right] I_\nu = j_\nu - \kappa_\nu I_\nu \tag{2.2.1}$$

where $I_\nu$, $j_\nu$, $\kappa_\nu$ are the monochromatic specific intensity, emission coefficient per unit mass and absorption coefficient (opacity) at
frequency $\nu$ and $\rho$ is the material density. In general, these quantities are functions of time, space and angle. Henceforth, this dependence will not be explicitly written down and these quantities shall be referred to simply as the intensity, the emission coefficient and the opacity or material opacity. $s$ is a coordinate in the spatial direction of travel.

Equation 2.2.1 has a very simple interpretation. It is essentially an equation of continuity. The expression in brackets on the LHS of 2.2.1 is just the total derivative along $s$ because $I_\nu$ is in general a function of both space and time. The RHS simply states that the rate of change of the intensity along $s$ (LHS) is the difference between the rate of emission of radiation, $j_\nu$ and the rate of absorption of radiation, $\kappa_\nu I_\nu$ along $s$. In the problem of radiation transport in laser plasma, the radiation will tend to change macroscopically on a hydrodynamic time scale. This is because the radiation time scale $\sim L/c \text{ or } 1/(\rho \kappa c)$ where $(\rho \kappa)^{-1}$ is an average photon mean free path, is small compared to the fluid time scale $\sim L/v$ where $v$ is a characteristic speed with which information (and thus change) can be transmitted i.e. $\sim$ the sound speed. In general, the scale of a laser plasma $L \gg$ the average photon mean free path $1/(\rho \kappa)$ so that the ratio of the radiation to fluid time scales is $O(v/c)$. Typically this is less than $10^{-3}$. Therefore, the radiation field will appear to adjust itself almost instantaneously to changes in the plasma conditions and can be regarded as being in a succession of quasi-steady states as the system evolves. Hence, on a hydrodynamics time scale, the first term in brackets i.e. the time derivative in equation 2.2.1 can be dropped. This is not true of course in a system which is evolving on a similar time scale to the radiation.
Equation 2.2.1 in general, does not have a simple solution. This is because $j_{\nu}$ and $\kappa_{\nu}$ are not independent of the radiation field so that the equation is no longer linear in $I_{\nu}$. However, in LTE $j_{\nu}$ and $\kappa_{\nu}$ are determined by the material properties and not the radiation (2.1.1). Introducing the source function $S_{\nu} = j_{\nu}/\kappa_{\nu}$ and optical depth $d\tau_{\nu} = \rho_\kappa ds$ (assuming $\rho_\kappa$ independent of $s$) equation 2.2.1 can be rearranged to give

$$\int_{I_{\nu}}^{I_{0\nu}} \frac{dI_{\nu}}{S_{\nu} - I_{\nu}} = \int_{\tau_{0\nu}}^{\tau_{\nu}} d\tau_{\nu}$$

where $I_{\nu}$ and $I_{0\nu}$ are the intensities at optical depths $\tau_{\nu}$ and $\tau_{0\nu}$ respectively. Performing the integration and using Kirchhoff's law $S_{\nu} = B_{\nu}$ one obtains

$$I_{\nu} = B_{\nu} \left[1 - \exp(-\Delta\tau_{\nu})\right] + I_{0\nu} \exp(-\Delta\tau_{\nu}) \quad 2.2.2$$

where $\Delta\tau_{\nu} = \tau_{\nu} - \tau_{0\nu}$. The second term on the RHS states that the contribution to $I_{\nu}$ due to $I_{0\nu}$ coming from $\Delta\tau_{\nu}$ away is just $I_{0\nu} \exp(-\Delta\tau_{\nu})$ and comes directly from the definition of the opacity. The first term on the RHS is the contribution to $I_{\nu}$ from emission in $\Delta\tau_{\nu}$. In the optically thick limit ($\Delta\tau_{\nu} \to \infty$) $I_{\nu} \to B_{\nu}$ irrespective of $I_{0\nu}$.

Equation 2.2.2 is the basis of the radiation transport model described in the next chapter. However, one further result will be required. It is useful to define the radiative flux $F_{\nu}$ which gives the net energy flow due to radiation of frequency $\nu$ passing through unit surface area in unit time. In an isotropic radiation field this is clearly zero and it is more instructive to define the energy flow through the surface into one particular hemisphere rather than the sum of the flows into both halves of the sphere about the surface. Because the field is isotropic, the only contribution to the energy flow will result from components of
I_v perpendicular to the surface, i.e. I_v cosθ because all the tangential components will exactly cancel. θ is the angle between I_v and the normal to the surface. Thus

\[ F_v^\uparrow = \int_{2\pi} I_v \cos \theta \, d\Omega = nI_v \]  \hspace{1cm} 2.2.3

where ↑ represents opposite directions, dΩ is an element of solid angle and the integration is performed over a hemisphere (2π). Therefore, a simple integration of equation 2.2.2 multiplied by \cos θ over a hemisphere will give the desired relation detailing how the radiative energy is being transported.

There is one final comment on equation 2.2.1. The opacity must be modified to take account of stimulated processes. It is easily shown that the correction factor in LTE is \((1 - \exp(-hv/kT))\) (Armstrong and Nicholls 1972). This factor applies only to true absorption and not scattering because stimulated scattering into and out of a beam of radiation exactly cancel (Mayer 1948). Therefore, the opacity that appears in equations 2.2.1 and 2 and in Kirchhoff's law must inherently include this correction factor.

### 2.3 Opacity Contributions

There are only three sources of opacity of any significance in laser plasmas. High energy processes such as nuclear absorption, pair production and inverse Compton scattering are completely negligible. Also, it was concluded in the first part of this chapter (2.1.3) that Compton scattering can also be ignored. This leaves f-f, b-f and b-b absorption.
f-f absorption (inverse bremsstrahlung) is the process by which a photon is absorbed by a free electron which then makes a transition to another free state. This process can only occur during an electron-ion collision so that energy and momentum may be conserved simultaneously, the majority of the photon's energy being taken up as electron kinetic energy. For a plasma at temperature $T$, consisting of ions of charge $Z^*$, the cross-section per ion for f-f absorption averaged over a Maxwellian velocity distribution assuming a Coulombic potential is given by (Clayton 1968)

$$\sigma_{\nu ff} = \frac{k_{ff} Z^{*2} n_e g_{ff}}{T^{1/2} \nu^3}$$

where $k_{ff}$ is a numerical constant, $n_e$ the electron density and $g_{ff}$ the f-f Gaunt factor. The Gaunt factor (see for example Karzas and Latter 1961) is the quantum mechanical correction to the semi-classical cross-section obtained by Kramers (1923).

A b-f transition occurs as a result of an electron initially bound to an ion, absorbing a photon and making a transition into the continuum. This is also known as photoionization. For an electron, initially with principal quantum number $n$, whose binding energy $E_n$ is determined by a Coulombic potential characterised by an effective charge $Z_n$, the b-f cross section is given by (Clayton 1968)

$$\sigma_{\nu bf} = k_{bf} \frac{Z_n^4 g_{bf}}{n^5 \nu^3}; \quad h\nu \gg E_n$$

$$\sigma_{\nu bf} = 0 \quad ; \quad h\nu < E_n$$

where $k_{bf}$ is a numerical constant and $g_{bf}$ is the b-f Gaunt factor.

b-b absorption is the process by which an electron, initially bound to
an ion, absorbs a photon and is excited to a higher lying bound quantum level. The cross-section for a transition from state \( i \) to \( f \) is given by

\[
\sigma_{\nu}^{bb} = k_{bb} \, f(i \rightarrow f) \Phi(\nu)
\]

where \( k_{bb} \) is a numerical constant, \( f(i \rightarrow f) \) is the oscillator strength for the transition and \( \Phi(\nu) \) is the line shape function which allows for the fact that the energy levels have a finite width.

The opacity can then be determined by summing all the transitions which occur at a particular frequency. For any particular transition it is given in general by

\[
\kappa_{\nu} = \frac{n \sigma_{\nu}}{\rho}
\]

where \( n \) is the number density of initial states. For f-f transitions, \( n \) is given by the electron density since the cross-section has already been averaged over all the momentum states. In the case of b-f and b-b transitions, \( n \) is given by the number of ions multiplied by the probability of the initial state existing.

The f-f and b-f cross-sections given above take no account of electron degeneracy. If the temperature is low enough and the density sufficiently high, then degeneracy can become important. In this case, the f-f cross section must be averaged over the Fermi-Dirac rather than Maxwell distribution function. In addition, because low lying momentum states will be (partially) filled, the probability that the final state is available, \( q(f) \), must be taken into account. The f-f and b-f cross-sections must be multiplied by

\[
q(f) = 1 - \frac{1}{\exp(E_f - \mu)/kT + 1}
\]
where \( E_f \) is the energy of the free electron involved in the transition, \( \mu \) the chemical potential and the second term on the RHS is the Fermi-Dirac occupation number. Hence, degeneracy has the effect of reducing the f-f and b-f opacities.

In the foil plasmas investigated in detail in chapter 5, degeneracy is not expected to be important for most of the simulation time. For example, in a carbon plasma at \( 1 \text{ gcm}^{-3} \) and 10 eV, \( \exp(-\mu/kT) \sim 1 \). Therefore, if \( E_f \sim 2kT \), i.e. 20 eV, \( q(f) \approx 0.9 \). Clearly, \( q(f) \) increases rapidly with increasing \( E_f \). Because most of the heating radiation from the gold exists above 100 eV and the target density is generally less than \( 1 \text{ gcm}^{-3} \), setting \( q(f) = 1 \) is a very good approximation for most of the radiation and is not expected to cause any significant errors in the results.
In this chapter, an experiment designed by the author to measure the transport of soft X-rays into the cooler, denser regions of laser targets using time resolved XUV spectroscopy is described. Both the experimental design, and diagnostics are discussed. A considerable improvement to the XUV spectrometer was made by the author in this experiment which was a follow up to a preliminary experiment. A highly polished, grazing incidence X-ray mirror, which acted as a high energy photon cut-off filter, was used to reflect the soft X-rays into the spectrometer, thus producing striking improvements to the spectra obtained, by the substantial reduction of higher spectral orders. These observations are the first time resolved measurements of radiation transport in this energy region through low Z materials. Both the experimental design and diagnostics are discussed before presenting some typical results, two of which are used in the extensive analysis using the radiation transport model described in the next chapter.

3.1 GENERAL DESCRIPTION

The experiment was carried out at the SERC Central Laser Facility (Rutherford Appleton Laboratory) using the Nd:glass laser Vulcan (Danson, Edwards and Wyatt 1985). Thin foil plastic targets of various composition and thickness (0.1 μm - 10 μm) were coated on one side with a thin layer (0.1 μm) of gold which was irradiated by a single beam of the laser. Frequency doubled, green (0.53 μm) laser pulses
700 ps (FWHM) in duration containing approximately 100 J were used. The laser light was focussed on target with an f/2.5 aspheric doublet lens to a focal spot around 400 μm (FWHM) in diameter producing an irradiance around $10^{14}$ Wcm$^{-2}$. The focal spot size was measured using an X-ray pinhole camera with a 50 μm resolution and filtered for photons in the 1.5 keV energy region. The soft X-rays produced in the hot gold plasma which were transmitted and transported through the plastic targets were analysed using a time resolved XUV spectrometer operating in the 10 Å to 70 Å spectral wavelength region. The temporal and spectral resolutions of the instrument were 50 ps and 250 mA respectively. A schematic of the experimental set up can be seen in figure 3.1.1.

3.2 TARGET DESCRIPTION

Most of the targets were coated on one side with a 0.1 μm thick layer of gold which acted as an intense radiation source. For comparison, some targets with no gold coating were also used. The targets were made of plastic, either N-parylene ($C_8H_8$), mylar ($C_{10}O_4H_8$) or formvar ($C_5H_7O_2$). The thickness of the targets varied between 0.1 μm and 10 μm. At 5 μm, all the targets were at least one photon mean free path thick, when cold, at all wavelengths within the detection region of the XUV instrument. The 0.1 μm substrate targets were specifically designed to be able to measure the emission from the rear of the gold since these substrates were too thin to affect the soft X-ray emission.

The targets were gold coated so that the source of X-rays was very intense, pseudo-continuous, approximately Planckian and could be
Figure 3.1.1. Schematic of experimental set up for radiation transport experiment.
considered localised in the gold plasma. Under the laser operating conditions, the 0.1 μm gold foil behaves ablatively and is expected to burn through only at the end of the laser pulse, if at all (see section 4.7.1). Therefore, the primary heating radiation is known to come from the gold i.e. the plastic was not heated directly by the laser.

3.3 DIAGNOSTICS

The X-ray pinhole camera is a standard piece of equipment. The time resolved XUV spectrometer has only been developed recently and is a powerful instrument. It consists of a flat field spectrometer, X-ray streak camera, image intensifier, and recording device (optical camera). Each of these individual devices is described in turn. Because the instrument is not absolutely calibrated, the relative instrumental response was estimated from details in the literature to be approximately constant in the operating range of the instrument. This is consistent with spectral measurements from ultra-thin 0.1 μm gold coated formvar substrates. The spectral profiles are well approximated by a Planckian spectrum with a temperature of 115 eV (to within a factor of approximately 2) in agreement with other measurements on gold targets (Mochizuki et al 1987).

3.3.1 Flat Field Spectrometer

The flat field spectrometer is certainly one of the most important XUV diagnostic developments to date for radiation transport studies. The most important element is the gold coated, grazing incidence, blazed grating (Harada and Kita 1980 and Nakano et al 1984) whose
main feature is that its focal plane is moved away from the Rowland circle and is flat over a large wavelength range. The flat focal field allows time resolved XUV spectroscopy to be performed simultaneously for a range of wavelengths using the same instrument. The flat field is achieved with variably spaced, curved grooves on a spherical surface. The grating used in this experiment has the following specifications: an average groove spacing of 1/2400 mm; a blaze angle of 1.9°; a 15.9 m radius of curvature and a 50\times30 \text{mm}^2 ruled area. The source to grating distance is fixed at 237 mm. A schematic of the grating operation is shown in figure 3.3.1. The distance $L_f$ is the distance from the grating centre to the focal plane and is a function of $\theta$, where $\pi/2 - \theta$ is the grazing incidence angle. This variation has been calculated by Nakano et al (1984) and is shown in figure 3.3.2 which has been taken directly from this paper. In the experiment, the grating was operated at a grazing incidence angle of 1.5° which corresponds to a focal length of 237 mm. From figure 3.2.2 it can be seen that the flat field can be expected to extend from around 10 A to above 100 A.

In the experiment, higher order contributions between about 10 A and 20 A were suppressed by the use of a highly polished silica mirror. The mirror was used at a grazing incidence angle of 2.5° to reflect the X-rays into the spectrometer. This acted as a high frequency cut-off filter for X-rays above about 1keV. Figure 3.3.3 shows the calculated reflectivity of the mirror as a function of photon energy for different grazing angles of incidence and is taken directly from Henke (1982).

Second order contributions from the grating were not clearly observed. Above approximately 30 A the signal level not in first order line
Figure 3.3.1. Schematic of operation of grazing incidence, blazed, flat field grating. $L_f$ is the distance from the flat focal field to the grating centre.
Figure 3.3.2. Calculated focal fields for different angles of incidence for a 2400 lines per millimetre, grazing incidence, blazed grating. The object to grating distance is fixed at 237 mm. Taken from Nakano et al (1984).
Figure 3.3.3. Calculated reflectivity of fused silica X-ray mirror versus photon wavelength for photons incident at a grazing angle of 2.5°. Taken from Henke et al (1982).
radiation was less than 5% of the measured first order signal in lines at lower wavelengths. Therefore, the second order contributions must be less than 5%. No higher order contributions have been observed at any wavelengths.

### 3.3.2 X-ray Streak Camera

The X-ray streak camera provides up to picosecond temporal resolution by first converting the photon signal into an electron signal via an appropriate transmission photocathode (in this case, low density caesium iodide, CsI). The secondary photoelectrons which are emitted with very low energies from the opposite side of the photocathode are then accelerated by a potential difference of several kV through a set of focussing electron optics. At the same time, the resulting electron beam is deflected by a ramp voltage of the order of 1 kV (depending on the sweep speed required) before it strikes a phosphorescent material at the opposite end of the instrument. Here the electron signal is converted into an optical signal which can then be amplified and recorded. Thus the signal observed at different positions on the phosphor along the direction of deflection corresponds to the photon signal striking the photocathode at different times. The relative separation of events is determined from the calibrated "streak rate" which is approximately constant across the extent of the phosphor and relates a spatial interval on the phosphor to a corresponding temporal interval.

A schematic of the X-ray streak camera and its operation is shown in figure 3.3.4. The temporal resolution of the instrument is determined by three factors for photons incident normally on the timing slit: the width of the timing slit; the spatial and energy spread of secondaries
Deflecting and Focussing Optics

Transmission Photocathode

Insulating Spacer

hv

Slit

Accelerating Grid

Anode

Phosphor

Temporal Direction

Figure 3.3.4. Schematic of X-ray streak camera operation.
from the photocathode and the resolution of the phosphor. It turns out that in the mode of operation with a timing slit of 500 μm and a streak rate of 100 ps mm⁻¹, the major factor determining the temporal resolution is the slit width. Because the timing slit is very close to the photocathode, the temporal resolution is effectively determined simply by the width of the slit. In this case, a slit width of 500 μm corresponds to a temporal interval and resolution of \((500 \times 10^{-3} \text{ mm}) \times (\text{magnification} = 1.1) \times (100 \text{ ps mm}^{-1}) \approx 50 \text{ ps}.

### 3.3.3 Image Intensifier and Optical Camera

Neither the image intensifier nor the optical camera introduce any variation in the relative response of the instrument with photon energy. The intensifier has six fixed available gain settings, each of which differs from the next by a gain factor of approximately 4. A more detailed discussion of the intensifier characteristics can be found elsewhere (Ridgeley 1987). For future reference, when the intensifier gain was set to minimum in the experiment, the relative gain is said to be 1. In general, a gain setting of \(n\) is said to have a relative gain (or gain) of \(4^n\). The optical camera uses Ilford Hp5 photographic film which is relatively calibrated using standard techniques.

### 3.4 EXPERIMENTAL RESULTS

To interpret the experimental data, it is also necessary to measure the dispersion of the grating in the focal field so that the wavelength at which particular spectral features occur can be identified. This could be estimated from figure 3.3.2. However, the reliability of the calculated curve is not known, and the instrument can only be set up
to within a particular tolerance each time (estimated to be \( \approx 2 \) A) and the streak camera and image intensifier introduce some slight distortion in any case. Therefore, the dispersion on the recording film is always measured from spectra for each particular set up. Figure 3.4.1 shows a streaked spectrum of a plastic containing nitrogen and oxygen with the instrument set in the configuration used in the experiment. Some of the easily observable lines are indicated above, while the dispersion curve determined from the line spectra of sodium, fluorine, oxygen, nitrogen and carbon is shown below. The dispersion varies from about 1.5 A per mm at 15 A to nearly 3 A per mm at 40 A.

3.4.1 XUV Streaked Spectra

XUV spectra were obtained for a considerable number of targets of different thickness and composition and for varying laser conditions. In the following, some typical examples of the streaked spectra obtained are presented. First, to emphasise the very significant improvement to the diagnostic by the addition of a grazing incidence X-ray mirror, spectra taken on 6 \( \mu \)m mylar targets from this and the preliminary experiment are shown in figures 3.4.2(a) and (b) respectively. Also shown are densitometer traces taken across wavelength at a time close to the peak emission.

The laser wavelength and pulse length were the same in both cases, 0.53 \( \mu \)m and 700 ps respectively. The spectrum from the first experiment was taken at a laser irradiance of \( 5 \times 10^{14} \) Wcm\(^{-2} \) while that from the second was taken at \( 1.5 \times 10^{14} \) Wcm\(^{-2} \). The difference in laser irradiance does not produce the distinct differences seen in figures 3.4.2(a) and (b) when the X-ray mirror is used. Note the absence of the sharp edge-like feature at about 45 A in the old data.
Figure 3.4.1 Streaked spectrum from a plastic target containing carbon, nitrogen and oxygen taken at a laser irradiance of approximately $10^{14}$ Wcm$^{-2}$. The instrument dispersion is also shown.
Figure 3.4.2(a). Streaked spectrum from the rear of a gold coated 6 μm mylar foil, obtained with the use of a grazing incidence X-ray mirror.

Figure 3.4.2(b). As figure 3.4.2(a) but without use of an X-ray mirror.
For future discussion, note the small shift in the edge feature towards higher energy (shorter wavelength) at later times in the spectrum from the new experiment. The apparently continuous spectrum produced in the old experiment is a result of multiple order contributions from the grating from radiation close to 1 keV. This is particularly apparent because the instrument filtering in the preliminary experiment (0.4 μm of aluminium on 0.4 μm of formvar) should have produced a significant (> factor of 10) and progressive reduction in the signal above 35 - 45 Å and a complete cut-off by 55 Å. This was not observed in the experiment because the signal in this region was dominated by higher order contributions. The instrument filtering in the second experiment was 0.1 μm of formvar which has negligible effect on the spectrum.

The longer emission time observed in the previous experiment is also largely a result of multiple order contributions, particularly because at the higher irradiance, the radiation in the kilovolt region is more readily produced. This can be seen in figures 3.4.3(a) and (b) which show spectra taken on 0.1 μm formvar targets and densitometer traces at a time close to the peak emission. Both were taken at similar irradiances to the respective values given above and at the same relative gain, 4. However, the filtering of the instrument was different for the two cases and was at least 1000 times more attenuating in the old experiment at around 60 Å. Despite this, the signal towards longer wavelength on this spectrum does not decrease accordingly, suggesting that it is composed of higher energy radiation in multiple orders. Direct laser heating of the plastic, occurring because the gold probably burns through before the end of the laser pulse at the higher irradiance, would result in a longer emission time although this
Figure 3.4.3(a). Streaked spectrum from the rear of a gold coated ultra-thin 0.1 μm formvar substrate, obtained without the use of a grazing incidence X-ray mirror.

Figure 3.4.3(b). As figure 3.4.3(a) but with use of an X-ray mirror.
emission would be much less intense (see figure 3.4.4).

All spectra shown henceforth were obtained in the experiment described in the previous sections. For comparison to figure 3.4.2 a spectrum taken at a laser irradiance of $1.7 \times 10^{14}$ Wcm$^{-2}$ from a 6 μm mylar target with no gold coating is shown in figure 3.4.4. A densitometer trace across wavelength close to the peak of the emission is also shown. This shot was taken with a relative gain of 256 in contrast to the gain of 16 for the spectrum shown in 3.4.2(b), while the filtering was the same in both cases. From the densitometer traces of the spectra from the 6 μm targets, it is clear that the emission from the plastic plasma is approximately 50 times less than that from the gold. Also, note the longer emission time for the plastic plasma. This is due to the lower temperature of the X-ray emitting region of the gold plasma because of the much more efficient radiative cooling.

In figures 3.4.5(a),(b) and (c) are shown streaked spectra and densitometer traces across wavelength close to peak emission from 3 μm, 5 μm and 10 μm N-parylene targets respectively. These data shots were taken at laser irradiances of $1.8 \times 10^{14}$ Wcm$^{-2}$, $1.5 \times 10^{14}$ Wcm$^{-2}$ and $1.7 \times 10^{14}$ Wcm$^{-2}$ respectively. The spectra from the 3 μm and 5 μm targets were taken with a relative gain of 16, while that from the 10 μm target was taken with a gain of 64. The filtering was 0.1 μm for the 5 μm and 10 μm targets and an extra 0.1 μm of formvar and aluminium for the 3 μm foil. On the spectrum from the 3 μm target, "pinching" of the spectrum can be observed at around 15 A and 45 A. This is due to overloading of the cathode because of a very high incident radiative flux. This can only
Figure 3.4.4. Streaked spectrum from the rear of a 6 μm mylar foil (no gold coating), obtained with the use of a grazing incidence X-ray mirror.
Figure 3.4.5(a). Streaked spectrum from the rear of a gold coated 3 μm N-parylene foil, obtained with the use of a grazing incidence X-ray mirror.

Figure 3.4.5(b). As figure 3.4.5(a) but for a 5 μm foil.
Figure 3.4.5(c). As figure 3.4.5(a) but for a 10 μm foil.
be prevented by filtering of the collected photon signal. However, filtering does introduce further undesirable non-linearities in the instrumental response. Note the increase in the delay in the emission of radiation just above (in energy) the edge feature observed in all three spectra at approximately 45 A with increasing target thickness. Also, the emission of radiation is progressively retarded towards longer wavelength for given target thickness. This is most emphasised in the spectrum from the 10 μm foil. As with the 6 μm, gold coated mylar target, the edge shifts towards higher energy at later times in the spectra from the 3 μm and 5 μm foils. Although there appears to be a slight shift in the edge in the spectrum from the 10 μm target, it is noticeably less than that observed from either the 3 μm or 5 μm foils.

3.4.2 Discussion
In all the results, except for those from the 0.1 μm substrates the most striking feature is the edge at 45 A. At later times, in the spectra from the thinner targets, the edge shifts towards higher energy by approximately 6 eV (~ 1 A). From figures 3.4.6(a) and (b), which show traces close to the peak emission across the spectra shown in 3.4.3(b) from a 0.1 μm substrate and 3.4.5(b) from a 5 μm foil, it is clear that this feature originates in the plastic foils and not in the gold plasma. The data presented in 3.4.5(b) is analysed in detail in chapter 5.

The position of the edge corresponds closely to the cold carbon K-edge position. An explanation based on this coincidence is not plausible, however, for the thinner targets (< 6 μm). As will be seen in chapter 5, a significant shift of the K-edge of more than 10 A towards shorter wavelength (higher energy) is predicted by the
Figure 3.4.6(a). Densitometer trace at peak emission across the spectrum shown in figure 3.4.3(b) from 0.1 \( \mu \text{m} \) of gold deposited on an ultra thin 0.1 \( \mu \text{m} \) formvar substrate.

Figure 3.4.6(b). Densitometer trace close to peak emission across the spectrum shown in figure 3.4.5(b) from 0.1 \( \mu \text{m} \) of gold deposited on a 5 \( \mu \text{m} \) N-parylene foil.
modelling in contrast to the measured shift of the order of 1 Å. It is believed that the presence of strong absorption lines is responsible for the discrepancy between the modelling and observations. This is discussed further in chapter 5.
CHAPTER 4

MULTIGROUP RADIATION TRANSPORT MODEL

In the following, the radiation transport and opacity model used to predict and analyse the results of an experiment detailed in chapter 3 are described. Initially, it is assumed that material opacities are known so that the transport model may be presented independently. The model uses an adaption of the multigroup formulation of the discrete $S_N$ method (Carlson 1963) in 1D. In this particular case the number of groups used is 116 and $N=2$. At first, the equations employed in the transport model are obtained for monochromatic quantities. These are then modified to the multigroup approximation. Considerable effort has been devoted to $S_N$ methods, a detailed discussion of which together with the various problems that can be encountered may be found elsewhere (Grant 1963, Lathrop 1968, Wendroff 1969, Lathrop and Carlson 1971).

The opacity code which employs an AA screened hydrogenic approximation in LTE to calculate group averaged Planck mean opacities at the material temperature is then described in detail. Screened hydrogenic potentials have been used because the opacities are calculated in-line with the hydrodynamics and speed of computation is essential. The subject of opacity calculations has been reviewed in a number of articles (Huebner 1964, Cox 1965, Carson Mayers and Stibbs 1968, Rozsnyai 1982, Rose 1984). Results from the opacity model are compared with measured data. Finally, it is explained how the model is applied to the dense foil plasmas.
experimentally investigated in chapter 3.

4.1 ADAPTATION OF $S_2$ FORMALISM

In the following, the approximation is made that the radiation energy density, pressure and momentum need not be considered in the hydrodynamics equation. This is because it was shown in 2.1.4 that these quantities contribute very little ($\lesssim 1\%$) to the total values. A general treatment of radiation hydrodynamics can be found elsewhere (Mihalas and Mihalas 1984). Therefore, to determine the energy redistribution due to radiation transport in the plasma it is necessary to calculate a net energy emission or absorption term. This must then be added to the hydrodynamics energy equation.

Although it is assumed that the plasma motion can be adequately described by a one dimensional hydrocode, it is necessary to take account of the fact that photons need not always, and in general do not, pass normally through each cell boundary. As a result, the distance travelled by a photon in a fluid cell between entering and exiting through opposite boundaries which are parallel, is not in general given by the perpendicular (shortest) distance between them. This is important because the attenuation of a beam of photons varies exponentially with the distance travelled in the medium. Therefore, on average, the energy deposition in a fluid cell is larger than if all the photons travelled normally to the cell boundaries.

Equation 2.2.1 will take account of this automatically because $s$ is a general spatial coordinate in the direction of the ray path. However,
it is necessary to relate s to the coordinates of the system. For a
general set of orthogonal curvilinear coordinates \( x^i \)
\[
\frac{\partial}{\partial s} = \frac{\partial x^i}{\partial s} \frac{\partial}{\partial x^i}
\]
4.1.1

The simplest 1D system to consider algebraically is a semi-infinite slab.
All quantities are then symmetric about an axis perpendicular to the
slab surface. The coordinate parallel to this axis will be denoted \( z \).
The radiation field can then be described by \( I_\nu(z, \theta) \) at any given time,
where \( \theta \) is the angle between the direction of \( I_\nu \) and \( z \). In this case,
\[
\frac{\partial x^i}{\partial s} = \cos \theta = \mu
\]
4.1.2

and equation 2.2.1 can be rewritten to give
\[
\rho^{-1} \left[ \frac{\partial}{\partial t} + \mu \frac{\partial}{\partial z} \right] I_\nu = J_\nu - \kappa I_\nu
\]
4.1.3

The discrete \( S_N \) form of this equation is given by (Hill 1971)
\[
\rho_{i - 1/2}^{n - 1/2} \left[ \frac{I_i^n - I_{i-1/2}^{n-1}}{c\Delta t} + \mu_{j-1/2} \frac{I_i^{n-1/2} - I_{i-1}^{n-1/2}}{\Delta z} \right] = j_{i-1/2}^{n-1/2} + \kappa_{i-1/2}^{n-1/2} I_{i-1/2}^{n-1/2}
\]
4.1.4

where the \( \nu \) subscripts have been suppressed for clarity. These
quantities are illustrated in figure 4.1.1 and are all in the discrete
direction \( \mu_{j - 1/2} \). The superscripts refer to the time step; \( n-1/2 \) is half
way between \( n-1 \) and \( n \). The \( i \) subscripts refer to the spatial mesh and
\( i-1/2 \) to the value at the cell centre. The \( j \) subscript refers to the
discrete direction and \( j-1/2 \) is the value half way between \( j-1 \) and \( j \).
All centred values \((-1/2)\) are found by simple averaging of the
boundary values.
Figure 4.1.1. Schematic of the quantities appearing in the 1D discrete $S_N$ form of equation 4.1.3 given by equation 4.1.4. $z$ is the spatial and $t$ the temporal direction.
It was shown in section 2.2 that at any instant in time, the first term in equation 4.1.4 can be ignored provided the radiation field is "frozen" to the plasma, i.e. radiation relaxation time scale \( \ll \) hydrodynamic time scale. In the particular case of the experiment described in chapter 3 this is expected to be true. If it is also assumed that any external radiation sources vary on a hydrodynamic time scale (e.g. radiation from the dense part of another laser plasma) then, provided retarded effects can be ignored, the time dependent term on the LHS of equation 4.1.4 can be dropped. The last condition is necessary for reasons of causality. It is clear that for any two points separated by a distance \( d \) at time \( t \), each will see radiation which was emitted from the other at time \( t - d/c \). Therefore, if a source emits radiation described by \( I_\nu(t) \), an observer a distance \( d \) away will measure radiation characterised by \( I_\nu(t - d/c) \) at time \( t \). If a time interval \( \Delta t_f \) is considered such that \( I_\nu(t + \Delta t_f) \approx I_\nu(t) \) and \( \Delta t_f \gg d/c \) then at any time \((t, t+\Delta t_f)\), \( I_\nu(t-d/c) \approx I_\nu(t) \). For the case of the foil plasmas described in chapter 3, retarded effects can be ignored, provided the light crossing time \( L/c \ll \Delta t_f \) the hydrodynamics time step. In general, \( \Delta t_f \approx 1 \) ps and \( L \approx 10 \mu m \) so that the above condition appears to be easily satisfied. Because the foils are continuously expanding, it is not obvious that this will always be the case. However, comparison of the hydrocode time step with the light crossing time of the target showed that, except in the initial stages, the latter was less than 5% of the former during most of the simulation. Therefore, the time dependent term in equation 4.1.4 can be ignored. Of course, in a strict sense, the other values of \( I_\nu \) in equation 4.1.4 should still be time centred. However, because \( \Delta t_f \) is small it is assumed that any quantity \( a^n - 1 \approx a^n \) so that \( a^{n - 1/2} \approx a^{n-1} \). Because the radiation energy density, pressure and momentum are not included in
the hydrodynamics equations, this last approximation allows the radiation transport calculation to be decoupled from the iterative solution for the hydrodynamics quantities, i.e. it can be performed at the end of the time step.

With these approximations equation 4.1.4 has been considerably simplified and is no longer required because equation 2.2.1 is readily integrated in these circumstances and in the special case of LTE considered here gives equation 2.2.2. Using relations 4.1.1-2, equation 2.2.2 can be rewritten

\[ I_\nu = I_{\nu 0} \exp(-\Delta z \rho \kappa_\nu / \mu) + B_\nu \left[ 1 - \exp(-\Delta z \rho \kappa_\nu / \mu) \right] \]

This equation is precisely the form required to transport the radiation through the hydrodynamics mesh if \( \Delta z \) is the distance between the cell boundaries and \( \rho, B_\nu \) and \( \kappa_\nu \) are the average values in the cells (i.e. values at cell centres). In discrete form the above equation may be written as

\[ I_i = I_{i-1} \exp \left[ -\frac{(z_i - z_{i-1})}{\mu_{j-1/2}} \rho_{i-1/2} \kappa_{i-1/2} \right] + \]

\[ B_{i-1/2} \left[ 1 - \exp \left[ -\frac{(z_i - z_{i-1})}{\mu_{j-1/2}} \rho_{i-1/2} \kappa_{i-1/2} \right] \right] \]

where the \( \nu \) subscript has again been suppressed. No time dependence is indicated because of the assumptions made above.

Although equation 4.1.5 is not a discrete \( S_N \) equation, with the approximations made, it could just as easily have been obtained from equation 4.1.4. This should be expected because without time dependence, equation 4.1.4 is the discrete \( S_N \) form of equation 2.2.1. Similar results are given by Lathrop (1969), Wendroff (1969) and
Hankin, Hill and Warham (1971) all of whom used a characteristics approach to obtain time dependent difference equations which reduce to equation 4.1.5 in the steady state. In general, provided the assumptions made above are accurate, equation 4.1.5 will be more accurate than the steady state $S_N$ equation. This is because equation 4.1.5 is an exact solution of the latter.

In order to determine the radiative flux at each cell boundary, the above equation must be summed over all $j$ (integrated over all solid angles). Using a bidirectional approximation in 1D, only $\mu > 0$ and $\mu < 0$ are distinguishable i.e. the forward and backward directions (Rose and Evans 1983). Also, assuming that $I_\nu$ is isotropic in each half plane ($\mu > 0$, $\mu < 0$) and replacing $(z_i - z_{i-1})/\mu_{i-1/2}$ by some mean chord $\langle \Delta r \rangle$, equation 4.1.5 can be multiplied by $\mu$ and integrated over each half plane separately with respect to solid angle (c.f. equation 2.2.3) to give

$$
\begin{align*}
F_i^+ &= F_i^+ \exp(-\langle \Delta r \rangle_k \rho_k K_k) + S_k \\
S_k &= \pi B_k \left[ 1 - \exp(-\langle \Delta r \rangle_k \rho_k K_k) \right]
\end{align*}
$$

4.1.6

where $k$ refers to cell values and $+$ and $-$ refer to the forward and backward direction respectively with $i$ decreasing inwards (forwards).

Hence, the seemingly very complicated problem represented by equation 4.1.4 has been reduced to the comparatively simple one represented by equation 4.1.6. However, application of equation 4.1.6 to any but planar geometries will rapidly lead to a violation of energy conservation and unphysical cell heating. This is not surprising because planar geometry was considered in deriving the
equation. It is easily seen by considering an empty cell (no sources or sinks of radiation) \( k \), in spherical geometry with \( i \) increasing outwards (figure 4.1.2). In the absence of sources and sinks, no energy can be deposited or generated in the cell and

\[
F_i^+ = F_{i+1}^+
\]

which is predicted by equation 4.1.6 and the energy passing through each boundary per unit time is given by \( F_i^+A_i \) and \( F_{i+1}^+A_{i+1} \) respectively. However, rays passing through point \( P \) on the outer boundary will only pass through the inner boundary if they are confined to the cone with solid angle \( \alpha \). The remainder of the energy travels back out through \( i+1 \) appearing as outwardly directed flux given by

\[
G_{i+1} = (F_{i+1}^+A_{i+1} - F_i^+A_i) / A_{i+1}
\]

For an absorbing cell, \( k \), in general, the contribution to \( F_i^+ \) due to radiation which entered through \( i+1 \) and is not absorbed in \( k \), is given by

\[
F_{i+1}^+ \exp(-<\Delta r>_k \rho_i k_k)
\]

Therefore, using this and the above expression for \( G_{i+1}^+ \), the radiation which enters through \( i+1 \) and is not absorbed in \( k \) but which exits as outwardly directed flux through \( i+1 \) is characterised by

\[
G_{i+1} = F_{i+1}^+ \exp(-<\Delta r>_k \rho_i k_k) \left( 1 - \frac{A_i}{A_{i+1}} \right)
\]

This quantity must be added to equation 4.1.6 when the outwardly directed flux is considered giving

\[
F_{i+1}^- = F_i^- \exp(-<\Delta r>_k \rho_i k_k) + S_k + G_{i+1}^-
\]

These equations are now quite general for any 1D geometry. Equation 4.1.6 is used for the inwardly directed flux and 4.1.7 and 4.1.8 are used for the outwardly directed radiation.
Figure 4.1.2. Diagram of cell k in spherical geometry showing that rays passing inwards through point P on the outer boundary which are confined to the solid angle $\Omega$, pass through the inner boundary, e.g. $r_1r_2$, while those outside the cone defined by $\Omega$ do not, e.g. $r_3r_4$. 
In planar geometry equation 4.1.8 reduces to 4.1.6 as expected since all the $G_i^r = 0$.

Finally, $<\Delta r>$ is approximated as the "mean" chord given by a formula due to Dirac as four times the volume divided by the total area (Armstrong and Nicholls 1972). For a cell in 1D this is given by

$$<\Delta r>_k = \frac{4V_k}{A_i + A_{i+1}}$$

4.2 MULTIGROUP REPRESENTATION

The quantities appearing in equations 4.1.6 - 4.1.8 are monochromatic quantities, the $\nu$ subscript being suppressed for brevity. In order to determine the energy transport through the system, these quantities must be integrated over the whole photon energy spectrum. Because the opacity is strongly frequency dependent, this cannot be done analytically, and numerical techniques must be used. The method used here is the multigroup technique in which the photon energy spectrum is divided into a number of discrete groups. The total radiative flux in each group is then defined as

$$F^g = \int_{\nu}^{\nu + \Delta\nu} F_\nu \, d\nu$$

where $\Delta\nu$ is the group width. This is equivalent to defining a "specific" flux $F^g$ which is constant over the group with $F^g = F^g\Delta\nu$. A suitably averaged group mean opacity $\kappa^g$ is then calculated and used in the transport equations.

In LTE, the group source functions are given by
\[ S^g = \int_{\nu}^{\nu+\Delta\nu} B_{\nu} d\nu = \alpha^g \int_{0}^{\infty} B_{\nu} d\nu = \frac{\alpha^g \sigma T^4}{\pi} \tag{4.2.1} \]

where \( \alpha^g \) is the fraction of the Planck function in group \( g \), \( \sigma \) is Stefan's constant and \( T \) is the absolute material temperature. The groups are then independent and can be transported separately. The transport equations 4.1.6-8 in group notation become

\[
\begin{align*}
F_{i+1}^g &= F_{i+1}^g \exp(-\langle \Delta r \rangle_k \rho_k \kappa_k^g) + S_k^g \\
S_k^g &= \pi S_k^g \left[ 1 - \exp(-\langle \Delta r \rangle_k \rho_k \kappa_k^g) \right] \\
G_{i+1}^g &= F_{i+1}^g \exp(-\langle \Delta r \rangle_k \rho_k \kappa_k^g) \left[ 1 - \frac{A_i}{A_{i+1}} \right] \\
F_{i+1}^g &= F_i^g \exp(-\langle \Delta r \rangle_k \rho_k \kappa_k^g) + S_k^g + G_{i+1}^g
\end{align*} \tag{4.2.2}
\]

where \( S_k^g \) is given by equation 4.2.1.

The energy deposited in each cell in time \( \Delta t \) is determined by the difference between the energy in each group entering the cell and that exiting the cell in \( \Delta t \) summed over all the groups

\[
\Delta E_k = \sum_g \left[ \left( P_{i+1}^g + P_{i}^g \right) - \left( P_{i}^g + P_{i+1}^g \right) \right] \Delta t \tag{4.2.3}
\]

where \( P_{i}^g = F_{i}^g A_i \). The energy flow per unit time is illustrated for spherical geometry in figure 4.2.1. Equations 4.2.1-3 represent the radiation transport model which is used in the code. It is applied at the end of each time step after the hydrodynamics has been advanced as follows. Beginning with a suitable boundary condition at the maximum value of \( i \), which is the end at which the laser is incident, (e.g. the outside of a microballoon) the inward transport is calculated
Figure 4.2.1. Diagram depicting the radiative energy $F^g A$ passing through the walls of cell $k$ in spherical geometry per unit time.
using the first two equations of 4.2.2. Then, beginning with a suitable inner boundary condition at \( i=1 \), the outward transport is calculated using the remaining equations in 4.2.2. The energy deposited in each cell is then calculated by equation 4.2.3 and is added to the electronic energy.

### 4.3 GROUP AVERAGED MEAN OPACITIES

It can be shown that, in the diffusion limit, the correct total radiant energy transport and momentum balance can be achieved by using a single mean opacity for the whole spectrum (Mihalas and Mihalas 1984). That is the Rosseland opacity given by

\[
\frac{1}{\kappa_R} = \frac{1}{\kappa} \int_0^\infty \frac{\partial B_\nu}{\partial T} \, d\nu
\]

\( \kappa_R \) is a harmonic mean and therefore particularly sensitive to windows (i.e. low values) in the opacity spectrum. Hence, an accurate description of the opacity between strong absorbing features such as lines is essential.

Alternatively, it can be shown that in LTE, in order to obtain the correct total emission, the Planck mean is required

\[
\kappa_P = \frac{\int_0^\infty B_\nu \, d\nu}{\int_0^\infty \nu B_\nu \, d\nu} \quad 4.3.1
\]

In addition, if the Planck opacity is used in the optically thin limit
then the correct emission rate from the plasma is achieved. However, in the optically thick or diffusion limit, $\kappa_p$ does not give the correct flux.

There are several reasons for choosing group averaged Planck mean opacities rather than the Rosseland equivalents. Firstly, the justification for the use of $\kappa_p$ in general lies in the zeroth moment of the transfer equation which in discrete multigroup form is represented by equations 4.1.6-8, whereas that for the Rosseland mean lies in the first moment of the transfer equation. Secondly, the diffusion approximation, upon which the validity of the Rosseland weighting relies, is not valid in laser plasmas. Thirdly, the Planck opacity is more sensitive to the high points in the opacity spectrum which are far more easy to calculate than for instance the opacity in the far wings of a line. Finally, the computation of $\kappa_p$ is significantly easier than that of $\kappa_R$. This is because $\kappa_p$ is a straight mean. This allows individual Planck means to be calculated for each photoabsorption process contributing to the total opacity and then added to give the total Planck means. Furthermore, it will be seen in 4.6 that the Planck mean can be calculated analytically to a high degree of accuracy. This is not the case for the Rosseland means. In any case, as the group widths become narrower, the difference between the differently weighted group mean opacities decreases and in this particular case, the group width is only 10 eV in the energy region where the majority of the radiation in laser plasmas is produced. The group averaged Planck mean, $\kappa^g$ is defined as in equation 4.3.1 except that both integrations are carried out over the group width and not the whole spectrum. The calculation of $\kappa^g$ will be described in the latter part of this chapter.
4.4 INTEGRATION OF THE PLANCK FUNCTION BETWEEN FINITE LIMITS.

The Planck function $B_\nu$ appears in the LTE transport equation as well as in the Planck mean opacity. The integral of $B_\nu$ over the entire energy spectrum can be performed analytically and gives the well known result

$$\int_0^\infty B_\nu \, d\nu = \frac{\sigma T^4}{\pi}$$

However, in the multigroup formulation, the integral must be carried out over a specific frequency interval. In this case, an exact analytic solution cannot be found. However, it is shown here that the integral can be calculated using a simple summation series representation. This was found to be considerably faster and more accurate than numerical integration techniques such as Simpson's method.

The Planck function in frequency space is given by

$$B_\nu = \frac{2 \ h}{c^2} \frac{\nu^3}{\exp(h\nu/kT) - 1}$$

Substituting $x = h\nu/kT$ and integrating between $x_1$ and $x_2$ gives

$$I = \frac{2 \ k \ T^4}{c^2 \ h^3} \left( \frac{x_2^3}{e^x - 1} \right)_{x_1}^{x_2}$$

The denominator in the integral may be written as

$$e^{-x} (1 - e^{-x})^{-1} = \sum_{n=1}^{\infty} e^{-nx}$$

Using this the integral in $I$ may be written
\[ I = \int_{x_1}^{x_2} e^{nx} \, dx \]

Repeated integration by parts and rearranging gives

\[ I = \left[ \sum_{n=1}^{\infty} \left( \frac{x^3}{n^3} + \frac{3x^2}{n^2} + \frac{6x}{n^3} + \frac{6}{n^4} \right) e^{-nx} \right]_{x_1}^{x_2} \]  

4.4.1

For \( x \gg 1 \), \( I \) is rapidly convergent because of the \( e^{-nx} \) factor. The number of members of the sum required for a given \( x \) to obtain a convergence of better than 1 in \( 10^{14} \) are tabulated in table 4.4.1. For the foil plasmas, described in chapters 3 and 5, the maximum temperatures are of the order of 100 eV, below which there are only 10 energy groups in the transport code. Therefore, for most of the groups for most of the time, equation 4.4.1 can be and is used in the code. However, Simpson's method is applied for \( I \) wherever \( x < 1 \) and is accurate to approximately 1 in \( 10^8 \).

Finally, \( \alpha^8 \) is given by

\[ \alpha^8 = \frac{I}{\pi^4/15} \]  

4.4.2

where \( I = \pi^4/15 \) when \( x_1 = 0 \) and \( x_2 = \infty \).
<table>
<thead>
<tr>
<th>a</th>
<th>b</th>
<th>n</th>
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<tr>
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<tr>
<td>1</td>
<td>2</td>
<td>27</td>
</tr>
</tbody>
</table>

Table 4.4.1. Number, \( n \), of terms of the summation series representation of the integral of the Planck function evaluated at \( x \)

\[ a < x < b \quad ; \quad x = \frac{hv}{kT} \]

required to obtain a convergence of better than 1 in \(10^{14}\).
4.5 **ONE ELECTRON SCREENED HYDROGENIC AVERAGE ATOM MODEL**

In the description of the radiation transport model, it was assumed that monochromatic opacities were already known. Here it is described in detail how these quantities have been calculated in-line with the hydrodynamics using an AA approximation in LTE. Because speed of calculation is essential for an in-line opacity calculation, screened Coulombic potentials are used. The model is based on the LTE version of XSNQ of Lokke and Grasberger (1977) and has been discussed in detail by Mayer (1948). Recently, More (1982) presented a screened hydrogenic atomic model derived from the WKB approximation. He found that for low degrees of ionization, the binding energies of the outer shells for high Z ions were predicted to be significantly higher by Mayer's model than by Hartree-Fock-Slater calculations of Scofield. However, here, good agreement was found between Mayer's model and calculations using an MCDF atomic physics code (Grant et al 1980) for ionization energies for low Z ions. Table 4.5.1 shows a comparison between ionization potentials calculated using the screened hydrogenic approximation using Mayer's screening coefficients and MCDF calculations for carbon. The agreement is best for ionization from inner shells. Therefore, while realistic results can be expected for low Z plasmas, it should be noted that application of this model to high Z plasmas will result in significantly underestimated degrees of ionization for moderate temperatures.
### Table 4.5.1. Ionization potentials calculated by an MCDF code and the screened Coulombic approximation.

<table>
<thead>
<tr>
<th>CONFIGURATION</th>
<th>IONIZATION POTENTIAL (eV)</th>
</tr>
</thead>
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<tr>
<td></td>
<td>MCDF</td>
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<td><strong>L- SHELL</strong></td>
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</table>
4.5.1 Initial Simplifying Approximation: Free Ions and Electrons

In reality, an isolated ion treatment for ionic wavefunctions is not strictly correct because of the electrostatic influence of all the neighbouring ions and electrons. An isolated ion approximation clearly becomes less appropriate as the plasma density increases. Similarly, a free electron treatment for electrons not considered bound is not strictly correct either. This approximation also becomes less adequate as the plasma density increases and as its temperature decreases. The degree of accuracy of these approximations for a particular electron is clearly influenced by the relative sizes of all the individual electrostatic interactions of the electron with all the other particles and the electron kinetic energy. If the kinetic energy is dominant then the electron's wavefunction is well approximated by that of a free electron (high temperature). On the other hand, bound electron wavefunctions will be well approximated by an atomic wavefunction of an isolated ion, provided the interaction of the electron with other nuclei and free electrons is negligible (low density). In the intermediate regimes where the electrostatic interactions are comparable to the electron kinetic energy but the electrons cannot be considered bound in the conventional sense, the problem becomes very complicated. Also, the distinction between a bound and free electron state depends solely on the relative sizes of the kinetic and electrostatic energy terms. Therefore, because the form of the terms remains essentially unchanged and only the relative sizes alter, it is clear that the transition from a bound to a free state (i.e. into the continuum) is far from abrupt as the electron energy is gradually increased. In reality, the atomic wavefunctions of neighbouring ions overlap and form smeared out energy bands near the continuum (cf energy bands in metallic solids). As the ions become closer, more tightly bound levels will become
affected and the smeared out bands will begin at lower quantum levels. Of course, this also has the effect of decreasing both the ionization potential and number of the remaining bound electrons. The ionization that occurs as a result of this is known as pressure ionization.

The treatment of this problem is extremely complex and beyond the scope of this work. A detailed discussion can be found elsewhere (Inglis and Teller 1963, Zimmerman and More 1980, More 1982). Here, the assumption is made that for most of the time very few electrons occupy these band states so that their contribution to the opacity is very small and can be ignored. This is equivalent to the assumption that the position of the bottom of the continuum is definite and only a specified number of bound states exist below it (see 4.7.3) while only free states exist above it. The justification for this is that the electrons in these bands must be only loosely "bound" to the ions precisely because the band is very close to the continuum. In addition, it is assumed that the correction to the atomic potentials due to the neighbouring ions and electrons can be represented on average by a spherically symmetric term in the wavefunction Hamiltonian. In this particular case, to preserve the analytic solution for the isolated ion eigenvalues, the correction is made after these have been determined. A more detailed discussion of this problem can be found in Rose (1983), where multicentre wavefunctions are discussed.

4.5.2 Model Equations And Justification

The discussion will be limited to a monatomic system for clarity in the equations. The generalisation to mixtures of elements is straightforward and will be described at the end of this section.
The equations representing the one electron AA model which are to be solved to obtain atomic level occupation probabilities and eigenvalues are

\[
f_n = \frac{1}{\exp(\epsilon_n/kT - \eta) + 1} = \frac{P_n}{2n^2} \quad 4.5.1
\]

\[
\epsilon_n - \Delta \epsilon_n = -0.0136 \frac{Z_n}{n^2} \left\{ \sum_{\kappa=1}^{n} \left[ 1 + \left( \frac{\alpha Z_n}{n - \kappa + \left( \kappa^2 - \alpha^2 Z_n^2 \right)^{1/2}} \right)^2 \right]^{-1/2} \right\} - 1
\]

\[
Z_n = Z - \left( \sum_m P_m \sigma_{n,m} - \frac{P_n \sigma_{n,n}}{2n^2} \right) \quad 4.5.3
\]

\[
Z^* = Z - \sum_n P_n \quad 4.5.4
\]

where all energies and temperatures (kT), are in keV. \( f_n \) is the occupation probability for level \( n \) which is determined from the Fermi-Dirac distribution function for independent fermions and is given by the number of electrons in the level, \( P_n \), divided by the level degeneracy, \( 2n^2 \). \( \epsilon_n \) is the \( n \text{th} \) level binding energy for the isolated ion. It is calculated using either the non-relativistic Bohr approximation (first expression in 4.5.2) or the solution to the Dirac equation, in which \( \alpha = 1/137.04 \) is the fine structure constant, averaged
over all the closely spaced levels $\epsilon_{n,\kappa}$ ($\kappa = 1, 2, \ldots, n$). $\kappa$ is a quantum number which replaces the orbital angular momentum number when the Dirac theory is used. $\Delta \epsilon_n$ is the correction to the isolated ion eigenvalues due to the neighbouring plasma and is discussed in more detail in 4.7.3. $\eta$ is a normalisation constant (chemical potential in units of kT) which is determined by the condition that

$$
n_e = \frac{4\pi (2 m_e)^{3/2}}{\hbar^3} \int_0^\infty \frac{\epsilon^{1/2} d\epsilon}{\exp(\epsilon/kT - \eta) + 1} = n_i Z^* \quad 4.5.5
$$

where $n_e$ and $n_i$ are the free electron and ion number densities respectively and the integral is known as the Fermi-Dirac function of order 1/2. The physical constants in front of the integral form the density of states for the free electrons. The calculation of the integral in equation 4.5.5 is described in 4.5.4. Here the high temperature approximation to $\exp(-\eta)$ which is used in the code is simply stated:

$$
\exp(\eta) = 3.154 \times 10^{-3} \frac{\rho Z^*}{A T^{3/2}} \quad 4.5.5(a)
$$

where $\rho Z^*/A$ is the electron density in mol cm$^{-3}$ and must take account of elemental mixtures (see 4.5.3).

$Z_n$ is the effective nuclear charge experienced by an electron in level n determined by 4.5.3 where $\alpha_{n,m}$ are screening constants (Mayer 1948). The second term on the RHS of 4.5.3 is the total screening from the nuclear charge for an electron in level n due to all electrons bound to the ion. The last term on the RHS of equation 4.5.3 is a correction to the sum which includes "self screening" because the sum is carried out over all electrons in all m including n. Finally, $Z^*$ is the average degree of ionization.
The independent electron treatment is justified by Mayer (1948). Briefly, he derived an expression for the number of electrons in a given energy region using a canonical ensemble treatment, including electronic interactions. He found that if terms up to first order in the electronic interaction energies only were retained, an equation of the form of 4.5.1 was readily obtained, provided $\epsilon_n$ is the energy of an electron in the $n^{\text{th}}$ level, taking into account electronic interactions. Therefore, an independent electron treatment is correct to first order, provided the electronic interactions are taken into account in the eigenvalues $\epsilon_n$ used in the expression for $f_n$. This is approximately achieved by using the screening constants $\sigma_{n,m}$ to calculate the effective nuclear charge experienced by an electron in level $n$. In fact, the $\sigma_{n,m}$ given by Mayer (1948) do not take account of the exchange interactions, which are comparatively small, but only of the spherically symmetric term in the Coulomb integral.

4.5.3 Generalisation to Elemental Mixtures and the Treatment of Hydrogen

Provided the ions in the plasma are not strongly coupled, it seems reasonable to assume that on average, all the ions will experience the same free electron density. Therefore, for each element in the mixture, equations 4.5.1-4 still apply to each individual element in turn but $e^{-n}$ must take into account the fact that $n_e$ is the same for each element. This is done by calculating $Z_i^*$ for each element $i$ using 3.5.4 and then calculating a mass averaged $Z^*$ to be used in 3.5.6. For example, suppose $A_i$ is the relative atomic mass of element $i$ which has an abundance by number of $X_i$ in the mixture, then the density is given by
\[ \rho = N_u \sum_i X_i A_i \frac{m_H}{\text{g cm}^{-3}} \]

where \( N_u \) is the number density of mixture "molecules" or "units" and \( m_H \) is the mass of a hydrogen atom. The number density of each element is given by the corresponding number of atoms of each element in a single mixture unit multiplied by the number density of units i.e.

\[ N_i = N_u X_i = \frac{X_i \rho}{m_H \sum X_i A_i} \text{ cm}^{-3} \]

The number density of electrons due to each element is then given by \( N_i Z_i^* \). In the case of hydrogen, the atoms are assumed to be fully ionized at all times. Because in general hydrogen only accounts for a very small proportion of the electrons in the plasma and does not significantly contribute to the opacity this approximation is not expected to incur any serious error. Therefore, the total number density of electrons is given by

\[ n_e = \frac{\rho \sum X_i Z_i^*}{\sum X_i A_i} \text{ mol cm}^{-3} \]

with \( Z_H^* = 1 \) always. For convenience this is written as \( \rho Z^*/A \).

4.5.4 Solving for the Atomic Level Populations

Equations 4.5.1-6 are solved using an iterative technique. Beginning with some initial set of \( P_n \), \( P_n^0 \), \( Z^* \) and the \( Z_n \) are calculated from equations 4.5.4 and 4.5.3. These are used to calculate the eigenvalues \( \epsilon_n \) from equation 4.5.2 and \( n \) from 4.5.5 (see 4.5.5) from which a new set of populations \( P_n^N \) are calculated. If convergence (defined later) is
not achieved, then a new set of populations \( P_n \) given by

\[
P_n = \begin{cases} 
(1 - a_c)P_n^O + a_cP_n^N & ; \ n = n_{\text{max}} \\
\frac{P_n^O}{P_n} & ; \ n \neq n_{\text{max}} 
\end{cases}
\]

is used in further cycles until convergence is obtained where, \( a_c \) is a damping factor used to inhibit numerical oscillations in the solution. Under some circumstances oscillations were found to prevent a solution (defined by satisfaction of the convergence condition) from being obtained. Oscillations are also suppressed by using a population \( P_n \) different from \( P_n^O \) only for the level \( n_{\text{max}} \) whose population changes by the largest amount for each element in the mixture at each iteration.

The convergence condition applied here is that the root mean square of the sum of the fractional changes in all of the populations (not just those that are altered for the next iteration cycle) be less than a pre-chosen fraction \( \epsilon \) for convergence to have been achieved i.e.

\[
\left[ \sum_i \sum_n \left( \frac{P_n^O - P_n^N}{P_n^O} \right)_i \right]^{1/2} \leq \epsilon
\]

Clearly, a larger value of \( \epsilon \) results in more rapid convergence but also gives a less accurate final solution. A value of \( \epsilon = 10^{-4} \) was found to produce very rapid results without sacrificing accuracy of the final solution. Almost identical results were obtained if \( \epsilon = 10^{-6} \) was used in transport calculations but took significantly longer to perform. Therefore, \( \epsilon = 10^{-4} \) was used in the code.

Convergence becomes more difficult to achieve as the temperature falls and the density rises. This is easily seen from equation 4.5.1. If
the temperature is sufficiently low that $\epsilon_n \gg kT$, then $P_n \sim 2n^2$. This eventually leads to an unphysical solution because the model will predict more bound electrons than there are protons in the nucleus. At this point, the model has broken down. This is related to the fact that the high temperature approximation to $\exp(-\eta)$, not valid in this plasma regime, is used in the Fermi-Dirac distribution function.

A valid solution of equation 4.5.1 relies upon the supposition that the chemical potential, $\mu$, can be accurately specified. For high temperatures, this is achieved by iteratively solving for $Z^*$, i.e. the number of free electrons. At low temperature, a further condition (automatically satisfied in the high temperature limit) is required that $Z^* > 0$ always. Physically, this occurs because the base of the continuum prevents all quantum levels with energies above a certain value existing as bound states thus restricting the number of bound electrons. Therefore, the accuracy to which $\mu$ may be calculated depends on the ability to determine accurately the position of the base of the continuum, i.e. the continuum lowering. This is a complicated problem and is not addressed here. Detailed discussions can be found elsewhere (Stewart and Pyatt 1966, More 1982, Burgess 1982).

If the plasma conditions are such that the model predicts a solution with $Z^* > 1$ for any element, then $Z^*$ is set to unity and the ions are filled with $Z - 1$ electrons using the auf-bau principle. This has the approximate effect of setting the opacities to their cold values. In this particular case, it is unlikely that densities sufficiently above solid occur to produce extreme pressure ionization or continuum lowering (see 4.7.3). The condition for a real solution is thus $Z^* > 1$. This has the added advantage that it prevents unphysical filling of upper shells.
independent of the maximum number of principal quantum levels at low temperatures. However, to prevent suppression of ionization particularly at low temperatures, the number of principal levels must be limited to a realistic value (see 4.7.3).

The number of iterations required for convergence not only depends upon $\delta$ but also on both the proximity of the $P^O_n$ to the final values and on the choice of $a_c$, the damping factor. The $P^O_n$ for each fluid cell in the laser hydrocode are taken from the values at the previous time step in the calculation. Because the conditions in the plasma in general change by only a small amount between successive time steps, this is probably the best possible guess and very close to the final solution for the actual populations.

The choice of $a_c$, as expected, was found to significantly affect the number of iterations required for convergence. In general, a larger value of $a_c$ provided a more rapid final solution. However, under certain conditions, particularly when the degree of ionization and temperature are low and the density is high, it was found that maximum values of $a_c$ exist above which solutions are not found due to the onset of numerical oscillations. Therefore, while a large value of $a_c$ is desirable, care must be exercised for the exact choice of $a_c$. This is illustrated in figure 4.5.1(a) and (b) in which optimised "$a_c$ lines" are shown for different densities against temperature for carbon and aluminium respectively. These were obtained by performing the described populations calculation at different densities and temperatures using values of $a_c$ between 0.1 and 1.0, three principal quantum levels and a $\delta = 10^{-4}$.
Figure 4.5.1(a). Optimised damping factors for iterative solution of average atom level populations against temperature for densities, $\rho$, of 0.1, 1.0, 5.0, 10.0 g cm$^{-3}$ for a carbon plasma. The vertical lines marked $T_c(\rho)$ show the temperature above which numerical oscillations never prevent convergence of $10^{-4}$ in the populations from being reached. Because only certain temperatures were sampled for each density, to ensure convergence below $T_c(\rho)$, a damping factor to the right of the shaded area for a particular density should be selected.
Figure 4.5.1(b). Same as 4.5.1(a) but for aluminium.
Each straight solid line results from interpolation between the optimised $a_c$ values determined, at the different temperatures sampled (1, 5, 10, 20, 40, 50, 75, 100, 200 eV) for a particular density. Also, shown for each density is a constant temperature line labelled $T_c(\rho)$ above which numerical oscillations never prevent convergence from being achieved, irrespective of the choice of $a_c$. Below $T_c(\rho)$ for a particular density, a choice of $a_c$ to the left of the solid line results in an oscillatory solution and no convergence for $\varepsilon = 10^{-4}$. Therefore, for convergence to be achieved, a value of $a_c$ to the right of the line should always be selected. Also, because $a_c$ was only optimised at selected temperatures for each density, the behaviour of the solutions between sampled temperatures, as $a_c$ is varied, is unknown. For this reason, a value of $a_c$ should always be chosen to the right of the shaded areas to ensure that the specified convergence of $10^{-4}$ is always achieved. As can be seen the $a_c$ lines are different for the two elements particularly below the $T_c(\rho)$ temperatures. Hence added care must be employed when choosing $a_c$ when elemental mixtures are considered.

The minimum value of $a_c$ shown is 0.1. In all the cases tested in figures 4.5.1(a) and (b) it was found that a convergence could be achieved at some temperature for which $0 < Z^* < 1$. Because the model prevents the equations from returning a solution with $Z^* < 1$, setting $a_c = 0.1$ never prevents (real) solutions with $Z^*$ only slightly greater than unity from being found.

4.5.5 Calculation of $e^{-\eta}$

The quantity $\exp(-\eta)$ is required to determine the level occupation probabilities given by equation 4.5.1 and is calculated from equation...
4.5.5. There is no general analytic solution to 4.5.5 and numerical techniques must be used. Perhaps the best of these and the most rapid is the method discussed by Cody and Thacher (1967). They used Chebychev polynomials to provide fits to the Fermi-Dirac function of order \( n \) (\( n = 1/2, 3/2 ... \)) accurate to about 1 in \( 10^8 \). However, the calculation of these is also time consuming and adds considerably to the complexity of the solution. Therefore, a very simple approximation to \( \exp(-\eta) \) is used in the model. It was found to be accurate to better than 4% and generally much less whenever a solution to the equations 4.5.1-5 could be found (i.e. \( Z^* \gg 1 \)). As noted earlier, the approximation becomes less valid at lower temperature and higher density when electron degeneracy becomes more important.

The approximation is as follows. If \( \eta \) is large and negative as in the non-degenerate case when the temperature is high and the density is low, then the denominator in equation 4.5.5 is given by \( \exp(-\eta) \exp(\epsilon/kT) \ (\gg 1 \text{ for all } \epsilon > 0) \). Therefore, the integral in equation 4.5.5 can be approximated by

\[
\exp(\eta) \int_0^\infty \exp(-\epsilon/kT) \, d\epsilon = \frac{h^3 n_e}{4\pi (2m_e)^{3/2}}
\]

Substituting \( u = (\epsilon/kT)^{1/2} \) reduces the integral to \( (kT)^{3/2} 2u^2 \exp(-u^2) \) which is easily integrated within the limits to give \( \pi^{1/2}(kT)^{3/2}/2 \). Replacing \( n_e \) with \( \rho Z^*/A m_H \) and evaluating all the constants one gets

\[
\exp(\eta) = 3.154 \times 10^{-3} \frac{\rho Z^*}{A (kT)^{3/2}}
\]

where \( kT \) is in keV and \( \rho \) is in g cm\(^{-3}\) and \( \rho Z^*/A \) is the electron density in mol cm\(^{-3}\) determined from equation 4.5.6.
Equation 4.5.9 is used in 4.5.1 to calculate the occupation probabilities. Figure 4.5.2 shows a plot of \( \exp(-\eta) \) calculated from 4.5.9 and also by numerical solution of 4.5.5 with a convergence factor of \( 10^{-6} \), against temperature for a carbon plasma at \( 1 \text{ g cm}^{-3} \). \( Z^* \) is calculated from equations 4.5.1-4 using 4.5.9 in place of 4.5.5. The \( Z^* \) calculated in this way were then used to calculate \( \exp(-\eta) \) numerically from 4.5.5. Also shown is the degree of ionization against temperature and the percentage difference between the \( \exp(-\eta) \) values calculated by the two different methods. Note that when \( Z^* > 1 \) (i.e. real solution of the equations) the percentage difference between the two methods is less than 4% and generally much less than 4% as the temperature increases. Therefore, whenever a real solution i.e. one with \( Z^* \neq 1 \) is found from equations 4.5.1-4 the expression 4.5.9 for \( \exp(\eta) \) is at most 4% in error but in general much more accurate. As the density is increased, the approximation given by 4.5.9 becomes less accurate. Equation 4.5.9 was used by Lokke and Grasberger (1977) and by Tsakiris and Eidmann (1987). However, neither commented on the accuracy of the approximation.

In the case of the foil plasmas described in chapters 3 and 5, the density is only predicted to rise just above \( 1 \text{ g cm}^{-3} \) during the initial stages of the simulations (4.7.3). Therefore, the approximation for \( \exp(\eta) \) given by equation 4.5.9 has been used in the model in place of equation 4.5.5.
Figure 4.5.2. Plot of degeneracy parameter against temperature, calculated for a 1 gcm$^{-3}$ carbon plasma from the high temperature approximation (---) given by the reciprocal of equation 4.5.9 using the degree of ionization, $Z^*$ (-----) calculated by the average atom model. Also shown is the degeneracy parameter calculated numerically (----) using the same electron density and the percentage difference, $\delta(D)$ (---) from the high temperature approximation.
4.6 CALCULATION OF GROUP AVERAGED PLANCK MEAN
OPACITIES FROM AVERAGE ATOM OCCUPATION NUMBERS

Only the b-f and f-f transitions are considered here. The shell
averaged Gaunt factors and f-f Gaunt factors are approximated by
unity. This then allows analytic expressions for the group averaged
Planck mean opacities to be derived. Equations 2.3.1-2 and 2.3.4 can
now be written in the following convenient form:

\[
\begin{align*}
\kappa_{\nu_{ff}} &= K_{ff} (\nu) \frac{3}{2} \\
K_{ff} &= 2.780 Z^* Z^* \rho \frac{1}{A} \frac{1}{(kT)^{1/2}}
\end{align*}
\]

\[4.6.1\]

\[
\kappa_{\nu_{bf}} = \sum_{i=1}^{n_{\text{max}}} \kappa_{\nu_{bf}(i)}
\]

\[4.6.2\]

\[
\kappa_{\nu_{bf(n)}} = K_{bf(n)} (\nu) \frac{3}{2}
\]

\[4.6.3\]

\[
K_{bf(n)} = 11.99 \frac{Q_n^4}{n^5} \frac{1}{A} P_n
\]

\[4.6.4\]

\[
Q_n = Z - \sum_{m=1}^{n_{\text{max}}} P_m \sigma_{n,m} + \frac{P_n \sigma_{n,n}}{2 n^2}
\]

where \(\nu\) and \(kT\) are in keV. \(Q_n\) is the inner screened nuclear charge
experienced by an electron in level \(n\). The justification for this
definition of \(Q_n\) can be found elsewhere (Slater 1960). In addition,
several different screening models for \(Q_n\) were assessed and that
described by equation 4.6.4 was found to give best agreement with
measured cold opacity data.
Following the discussion of 4.3, each photoabsorption process is considered separately. Using the multigroup form for the definition of the Planck mean, $\kappa^g$, (i.e. multigroup form of equation 4.3.1), the correction factor for stimulated emission given in 2.2 and equations 4.6.1-3, one obtains for the group averaged Planck mean

$$
\kappa^g_{\text{ff}, b(n)} = \frac{K_{\text{ff}, b(n)}}{(kT)^3 I} \int_{x_1}^{x_2} \frac{1 - e^{-x}}{e^x - 1} dx
$$

where $x = \frac{\hbar v}{kT}$, $x_1$ and $x_2$ are the group boundaries and it has been assumed that the material and radiation temperatures are equal. $I$ is defined in equation 4.4.1. The numerator in the integral is just the stimulated emission correction factor and the denominator arises from the Planck function. The $v^3$ dependence in equations 4.6.1 and 3 exactly cancel with that in the Planck function. Hence, the integrand reduces to $e^{-x}$ and is easily integrated to give

$$
\kappa^g_{\text{ff}, b(n)} = \frac{K_{\text{ff}, b(n)}}{(kT)^3 I} \left[ e^{-x} \right]_{x_1}^{x_2} 4.6.5
$$

Equation 4.6.5 is valid for all groups provided the opacity is continuous. Where, for example, a photoionization edge occurs in the middle of a group, the lower limit in the integral is simply replaced by $(\hbar v)_{\text{edge}}/kT$. However, $I$ is still calculated over the whole group as before.

To obtain 4.6.5 the assumption was made that the material and radiation temperatures were equal. This is correct when the emission from a body in LTE is considered because the emission coefficient is related to the opacity via Kirchhoff's law modified for stimulated
emission (2.2), provided the radiation temperature in the Planck function is equal to the material temperature. On the other hand, if an external radiation field is incident on the body, then 4.6.5 will not in general give the correct absorption of this radiation unless it is Planckian and also characterised by a temperature equal to the material temperature. However, because, in this case, the group width is very small (10 eV) in the region where most of the radiation exists, using equation 4.6.5 both for emission and absorption is not expected to cause any serious error. The total group averaged Planck means are then obtained by summing all the means for the individual processes occurring in each group. It should be noted that although setting the radiation temperature equal to the material temperature in the expression for the Planck means does not, in general, give exactly the correct absorption, neither would using two different temperatures (unless of course the external radiation field is exactly Planckian and has no effect on the level populations).

Equation 4.6.5 agrees very well with measured cold data for photoionization from inner shells. However, it predicts a very low opacity resulting from photoionization from outer shells. This is because the potential in the outer regions of the ion, particularly for low degrees of ionization, is far from Coulombic. The result is that the cross section falls off as $\nu^n$, $2 \leq n \leq 3$ rather than $\nu^3$ while the strength is still given approximately by equations 4.6.3-4 at the photoionization edge. In order to account for this, the frequency dependence of the b-f opacity for outer shells was set to $(h\nu)^{-1}_{edge}(h\nu)^{-2}$ for $h\nu \geq (h\nu)_{edge}$ while maintaining a Gaunt factor of unity. Equation 4.6.3 then becomes for outer shells

$$\kappa_{\nu bf(n)} = K_{bf} (h\nu)^{-1}_{edge} (h\nu)^{-2}$$

4.6.6
where $K_{b_l}(n)$ remains unchanged.

Equation 3.6.5 then becomes

$$
K_{\bar{n}, b_l(n)} = \frac{K_{\bar{n}, b_l(n)}}{(kT)^3 I} \left[ e^{-X} (1 + x) \right]^{x_1}_{x_2}
$$

4.6.7

The outer shells are defined as the valence shell of the cold atom and all levels above this. Therefore, for carbon and oxygen for example, all principal quantum levels outside and including the L-shell qualify as outer shells. Figure 4.6.1 shows the excellent agreement between measured L-shell opacity for cold mylar ($C_{10}O_4H_8$) (taken from Henke and Elgin 1970) below the carbon K-edge and that calculated by the model using equation 4.6.7 for the L-shell photoionization opacity (the f-f opacity is also included). Also shown are the results calculated from the model using equation 4.6.5 which lie significantly below the measured data. Equation 4.6.7 is used at all times for all outer shell photoionization. However, the radiation transport has been compared using both 4.6.5 and 4.6.7. The results are discussed in chapter 5. The overall agreement between calculated and measured values (taken from the same source) also appears to be excellent. Figure 4.6.2 shows the calculated and measured opacity (taken from the same source) for cold mylar in the 0.2 to 1 keV energy region.
Figure 4.6.1. Comparison between measured opacity of cold mylar (+) 
\((C_{10}O_{4}H_{10})\) below the carbon K-edge and that calculated using two 
different opacity scaling laws: \(1/(\hbar\nu)^2\) (---) and \(1/(\hbar\nu)^3\) (—).

(+) taken from Henke and Elgin (1970).
Figure 4.6.2. Comparison between measure opacity (+) of cold mylar (C_{10}O_{4}H_{8}) and values calculated by the author’s model (—) in the 0.2 - 1.0 keV energy region.

(+) taken from Henke and Elgin (1970).
4.7 APPLICATION TO EXPERIMENTAL CONDITIONS

First, it is explained how the numerical model is applied to the gold coated foil targets described in chapter 3. There are two points which are discussed: (1), the approximation of the behaviour of the gold layer (because it cannot be accurately calculated by the model); (2), the treatment of pressure ionization and continuum lowering.

4.7.1 Behaviour of the Gold Layer

Modelling of intense laser light interaction with high Z targets, in this case gold, is an extremely complex problem. Because, in general, around half of the absorbed laser energy is converted into soft X-rays, (Mead et al 1981, Kodama et al 1986, Nishimura et al 1983) it is clearly necessary to include an accurate radiation production and transport calculation in the model for realistic hydrodynamics and radiation calculations. To do this, one must account for both the marked non-LTE behaviour of the X-ray emitting region in the gold plasma and the large number of closely spaced lines. Even having successfully accomplished this, the radiation must then be transported. This problem is far beyond the scope of this work and could not, in any case, be accurately calculated by our model. Therefore, the behaviour of the gold is not calculated theoretically in the simulations. Instead, a time dependent X-ray spectrum, described in 4.7.2, is injected into one end (boundary) of a plastic foil of appropriate thickness and composition. This assumes that the majority of the plastic heating is due to X-rays from the gold. Other possible heating mechanisms include hot electrons, thermal electrons and laser heating. For most of the pulse, at least, laser heating can be eliminated. Mass ablation rate measurements for gold foils suggest
that with a green Gaussian laser pulse, 500 ps FWHM, only 600 A of gold are ablated at an irradiance of $10^{14}$ Wcm$^{-2}$ (Nishimura et al 1981, 1983). This gives a mass ablation rate of $4 \times 10^5$ gcm$^{-2}$s$^{-1}$ at a constant irradiance of $10^{14}$ Wcm$^{-2}$ which is in excellent agreement with other measurements on different materials (Goldsack et al 1982). Using this value and assuming that the 700 ps FWHM, $1.5 \times 10^{14}$ Wcm$^{-2}$ green laser pulse used in the experiment had a shape between a Gaussian and top hat, approximately 0.075 µm – 0.15 µm of gold can be expected to be ablated during the laser pulse. Therefore, burn through of the gold layer is only expected towards the end of the pulse, if at all.

Hot electron heating is assessed from measurements on gold disk targets illuminated by green laser pulses 600 – 700 ps FWHM at an irradiance of $3 \times 10^{14}$ Wcm$^{-2}$. They indicate a hot electron conversion efficiency of approximately 0.3% (Mead et al 1981) which is qualitatively in agreement with other measurements (Eidmann et al 1986, Mochizuki et al 1986, Sakabe et al 1988). Therefore, because the experiment described in chapter 3 was conducted at an even lower irradiance, heating due to suprathermal electrons is ignored.

The importance of foil heating due to electron thermal transport is estimated by comparing the radiation and electron thermal conductivities for the conditions expected to exist in the X-ray emitting region of the gold plasma ($\rho \sim 0.1$ gcm$^{-3}$, $T \sim 100$ eV (Mead et al 1983)). The coefficients are given by (Spitzer 1956, Zel’dovich and Raizer 1968, Mihalas and Mihalas 1984)

$$k_e \approx \frac{10^{-9} T^{5/2}}{Z^* \ln \Lambda} \text{ Jm}^{-1}\text{s}^{-1}\text{K}^{-1}$$
\[ k_R = \frac{3 \times 10^{-7} T^3}{(\rho \kappa_R)} \text{ Jm}^{-1} \text{s}^{-1} \text{K}^{-1} \]

where \( \ln \Lambda \) is the Coulomb logarithm and \((\rho \kappa_R)^{-1}\) is the Rosseland mfp in metre. \( \kappa_R \) is estimated from a power law given by Tsakiris and Eidmann (1987):

\[ \kappa_R = a T^s (\text{keV}) \rho^r (\text{gcm}^{-3}) \text{ cm}^{2g^{-1}} \]

For gold, they calculate \( a = 280, s = -1.058, r = 0.001 \). Therefore, \( \rho^s \approx 1 \) for any density encountered in the laser heated gold and \( \kappa_R \) is approximately linearly dependent on \( T \). Inserting the constants and assuming \( T = 100 \text{ eV} \) and \( \rho = 0.1 \text{ gcm}^{-3} \) and using \( Z^* \ln \Lambda \sim 100 \) \((Z^* \sim 20 - 30, \ln \Lambda \sim 5)\) one obtains

\[ \frac{k_e}{k_R} \approx \frac{3 \times 10^{-3} \rho \kappa_R}{Z^* T^{1/2} \ln \Lambda} \sim 10^{-3} \]

\( \ln \Lambda = \ln(D_{\text{max}}/D_{\text{min}}) \) was calculated assuming \( D_{\text{max}} = R_0 \), the ion sphere radius and \( D_{\text{min}} \) is the distance of closest approach obtained assuming a \( Z^* = 20 - 30 \) and an average electron thermal velocity.

Considering the above three points, it is expected that the majority of the foil heating is due to soft X-rays.

To approximate the effect of the pressure of the gold plasma in the hydrodynamics, the plastic boundary through which the soft X-rays are incident is prevented from expanding towards the "laser" and is held fixed throughout the simulation at its initial position. The validity and effects of these approximations is a source for further possible theoretical and experimental investigation.
4.7.2 X-ray Spectrum from the Rear of a Thin Gold Layer

In the simulations, the soft X-ray emission from the gold layer is calculated by using an analytic fit to the experimentally measured temporal variation at each wavelength with an assumed spectral profile at each time. This is because, although "monochromatic" temporal profiles can be obtained from spectra from ultra-thin substrate targets, the instrument is not absolutely calibrated and does not cover the entire wavelength region in any case. Therefore, the spectral profile is estimated from other measurements. These show the spectrum to be approximately Planckian with added features due to M, N and O-shell emission bands, for a wide range of laser conditions and target geometries (Herrmann et al 1986, Mochizuki et al 1986). Although these are all time integrated measurements, it is argued here that the spectrum near the peak of the radiation pulse must closely resemble the time integrated data because most of the emission occurs at that time.

Before estimating the total power in the radiation pulse, the temporal profile is determined because this is also required to calculate the total X-ray energy emitted from the gold, i.e the X-ray conversion efficiency. The temporal profile of the emission at each wavelength is obtained from the spectrum taken on the 0.1 μm formvar substrate target presented in figure 3.4.3(b). The FWHMs of several traces at various wavelengths were found to be approximately constant and equal to the nominal 700 ps FWHM of the laser pulse. Also, the relative time at which the leading edges reached half maximum was found to be approximately the same for each trace to within the temporal resolution (50 ps) of the streak camera. Therefore, a single approximate analytic form, f(t) ∝ f(ν,t), is used to describe the whole
spectrum. Above the half height, \( f(t) \) is given by a double Gaussian with the peak at \( t_0 \) and with \( 1/e \) times of \( \tau_L \) and \( \tau_R \) for \( t < t_0 \) and \( t > t_0 \) respectively. On the rising edge, below the half height, \( f(t) \) is joined smoothly to the Gaussian and extrapolated linearly to zero. \( f(t) \) is cut off after about two rise times. In normalised units, \( f(t) \) is given by

\[
\begin{align*}
  f(t) &= \frac{(\ln 2)^{1/2} t}{\tau_L} \quad ; \quad t \leq \frac{\tau_L}{2(\ln 2)^{1/2}}, \text{ i.e. } f(t) \leq 0.5 \\
  f(t) &= \exp \left\{ -\left[ \frac{(t - t_0)}{\tau_L} \right]^2 \right\} \quad ; \quad \frac{\tau_L}{2(\ln 2)^{1/2}} < t \leq t_0 \\
  f(t) &= \exp \left\{ -\left[ \frac{(t - t_0)}{\tau_R} \right]^2 \right\} \quad ; \quad t_0 < t \leq 2t_0 \\
  f(t) &= 0 \quad ; \quad t \geq 2t_0
\end{align*}
\]

4.7.1

Equation 4.7.1 is fit to the data with \( \tau_L = 516 \) ps and \( \tau_R = 324 \) ps. Figures 4.7.1 (a) - (e) show the comparison between \( f(t) \) and the five traces taken from the data shown in figure 3.4.3(b) across time at five wavelengths (17, 23, 31, 41, 51 A) in normalised intensity units. During the main heating part of the pulse, the approximate \( f(t) \) value is at most 25% different from the measured value. However, much better agreement ± 5% is obtained around 300 eV where most of the radiative energy is emitted. Simulations were carried out with slightly different rise times but this did not affect the results or conclusions drawn later. It is noted here that variation of the laser pulse shape and length will alter \( f(\nu,t) \). This is discussed in 5.5.2.
Figure 4.7.1(a) - (e). Comparison between approximate analytic fit (---) to the measured (—) temporal profile at five wavelengths (17, 23, 31, 41, 51 Å) from the rear of an ultra thin 0.1 μm formvar gold coated substrate irradiated by a $1.5 \times 10^{14}$ W cm$^{-2}$, 700 ps (FWHM) green laser pulse.
Figure 4.7.1(e).
This choice of f(t) allows the total energy delivered by the radiation pulse to be determined simply, given the total radiative power at the peak of the pulse because f(t) is wavelength independent. The conversion efficiency of laser to X-ray energy towards the rear of the target (i.e. the fraction of the laser energy that passes into the plastic in the form of radiation) is estimated from absolute time resolved diode measurements (Rose 1988). This data was taken with very similar laser conditions at an irradiance of $10^{14}$ Wcm$^{-2}$ from ultra-thin targets with an identical gold thickness. The total radiative emission from the rear of the foils was characterised by a temperature $T_R$ such that the total power emitted per unit area was given by $\sigma T_R^4(t)$. If $T_{RMAX}$ is the maximum measured temperature and $\tau_R$ and $\tau_L$ are the FWHM of the radiation and laser pulses respectively, then the conversion efficiency towards the rear is given by

$$\frac{\sigma T_{RMAX}^4 \tau_R}{I_L \tau_L}$$

where $I_L$ is the laser irradiance. In the experiment described in chapter 3, $\tau_L \approx \tau_R \approx 700$ ps. In the diode experiment it was found that $\tau_L \approx \tau_R \approx 350$ ps. Therefore, provided there is no strong dependence of the conversion efficiency on the laser pulse width or irradiance (Mochizuki et al 1986) $T_{RMAX}$ can be calculated for the experiment described here using

$$T_{RMAX} = I_L^{1/4} T_{RMAX(DIODE)}$$

where $I_L$ is the laser irradiance in units of $10^{14}$ Wcm$^{-2}$ and $T_{RMAX(DIODE)}$ was found to be 105 eV. For the data presented in figure 3.4.5(b) $T_{RMAX} = 115$ eV corresponding to an X-ray conversion efficiency of 12%. This value is in good agreement with other measurements (Mochizuki et al 1987). $T_{RMAX}$ will often be referred
to in place of the conversion efficiency. Figure 4.7.2 shows the measured spectrum at the time of peak emission from the data presented in figure 3.4.3(b) in arbitrary units. The band results from experimental uncertainties. Also shown are Planckian spectra with temperatures of 100, 115 and 130 eV which are normalised to the experimental curve at the wavelengths of peak emission, approximately 24, 22 and 19 Å. Although the absolute values of the experimental curve are not known, the Planckian spectrum with a temperature of 115 eV replicates the measured profile well (to within approximately a factor of 2) in this energy region. This is consistent with other measurements (Mochizuki et al 1987).

These three approximations now allow the total radiative flux entering the plastic at time t to be estimated. For this, it is assumed that the radiation be isotropic in the half plane because the ratio of the focal spot diameter to target thickness is very large. Therefore, the total radiative flux estimated to enter the plastic foils at time t is then taken to be

\[ F(t) = \sigma T_{R_{MAX}}^4 f(t) \]  \hspace{1cm} 4.7.3

where \( T_{R_{MAX}} \) is given by equation 4.7.2 and \( f(t) \) is given by equation 4.7.1. To obtain the monochromatic flux at each time t, it is assumed that, in the absence of any other data, the M, N and O-shell emission bands vary temporally in the same way as the remainder of the spectrum i.e. as \( f(t) \). If the fraction of the total radiative energy emitted in each of these bands above the continuum is given by \( \delta_M, \delta_N \) and \( \delta_O \) respectively, then the underlying continuum is described by

\[
\begin{align*}
F_{\nu}(t) &= \pi B_{\nu}(T_c) f(t) \\
T_c &= T_{R_{MAX}} \left[1 - \delta_M - \delta_N - \delta_O \right]^{1/4}
\end{align*}
\]  \hspace{1cm} 4.7.4
Figure 4.7.2. Plot showing the measured spectral profile from a streaked spectrum at time of peak emission from the rear of 0.1 μm of gold irradiated by a $1.5 \times 10^{14} \text{ Wcm}^{-2}$, 700 ps (FWHM) green laser pulse and supported on an ultra thin 0.1 μm formvar substrate. The shaded band results from experimental uncertainties. Also shown are Planckian profiles with temperatures of 100 (---) 115 (―) and 130 eV (----) which are normalised to the experimental curve at the wavelengths of maximum emission, approximately 24, 22 and 19 Å respectively.

* This comparison is not the primary justification for the use of $T_{\text{RMAX}} = 115 \text{ eV}$, which is estimated from other absolute measurements (see 4.7.2 p106).
and the additional flux due to the M, N and O bands is given by

$$F_V(t)_{M,N,O} = \sigma T_{RMAX}^4 \delta_{M,N,O} L_{\nu M,N,O} f(t) \quad 4.7.5$$

where $L_{\nu M,N,O}$ is a "band shape function" for the M, N and O band respectively such that

$$\int_{M,N,O} L_{\nu M,N,O} d\nu = 1 \quad 4.7.6$$

The group fluxes are then determined as in section 4.2.

With these definitions, it is always ensured that the total emission is given by 4.7.3. The effects of the inclusion of the various band emissions will be discussed in section 5.5.1.

4.7.3 Pressure Ionization and Continuum Lowering

Continuum lowering CL (to which pressure ionization (PI) contributes) cannot be attributed to any one individual mechanism and in dense plasma is an extremely difficult quantity to calculate (Zimmerman and More 1980, More 1982, Burgess 1982). When the plasma density is very high, CL and PI can reduce the number of quantum levels bound to an ion to just a few. In accordance, the number of principal quantum levels, N, used in the AA model must also be limited. This is particularly so for low temperature and high density, when the occupation probabilities of all the levels predicted by the model tend to unity. Clearly, if N is too large, the ionization of the atoms will be suppressed because of the $n^2$ dependence of the level degeneracies, i.e. electrons will be preferentially transferred to high n as opposed to the continuum (see figures 4.7.3 (a), (b) and (c)).
Figure 4.7.3(a). Comparison between degree of ionization versus temperature calculated using the screened hydrogenic average atom model with 2 (□), 3 (○) and 4 (+) principal quantum levels and an average atom model using a Thomas-Fermi potential (▲) for a carbon plasma at a density of 1 gcm⁻³.
Figure 4.7.3(b). Same as 4.7.3(a) but for a density of 0.1 gcm$^{-3}$. 
Figure 4.7.3(c). Same as 4.7.3(a) but for a density of 0.01 g cm$^{-3}$.
More (1982) discusses three different pictures for PI of a particular level. He shows that they each lead to the same result for the density at which PI of the level could be expected to occur. One argument is that the onset of PI of a level ought to occur when the ions become sufficiently close that the orbitals of the level begin to overlap. This occurs when the orbit radius, \( r_n \), is equal to the ion sphere radius, \( R_0 \). Approximating the \( r_n \) by a Bohr orbit leads to

\[
 r_n \approx 5.29 \times 10^{-11} \frac{n^2}{Q_n^2} \approx R_0 = \left[ \frac{3A m_H}{4\pi \rho} \right]^{1/3}
\]

where \( Q_n \) is the screened nuclear charge seen by the electron (see for example 4.5). Rearranging and evaluating the constants gives the density for the onset of PI

\[
 \rho(n) \sim 2.7 \frac{AQ_n^6}{n^6} \text{ gcm}^{-3}
\]

The range of conditions in the foil can be estimated as follows. If it is assumed that the foil expands isothermally at around the maximum temperature \( \sim 100 \text{ eV} \), then after 1 ns, a decompression of approximately 20 times for a 5 \( \mu \text{m} \) target results. Therefore, it can be expected that conditions in the foils vary from cold (1 gcm\(^{-3} \)) to approximately 100 eV and 0.05 gcm\(^{-3} \). In a cold carbon (A = 12) plasma, \( Q_n \sim 1 - 2 \) for \( n \gg 2 \) so that using the above expression

\[
 \rho(2) \sim 0.5 - 30 \text{ gcm}^{-3}, \quad \rho(3) \sim 0.05 - 6 \text{ gcm}^{-3} \quad \text{and} \quad \rho(4) \ll 0.05 \text{ gcm}^{-3}.
\]

Therefore, in the cold material, approximately the first 2 – 3 levels are expected to be little affected by pressure ionization. As the temperature rises \( Q_n \) increases, which suggests that higher \( n \) values are allowed before overlapping and thus PI of orbitals occurs. However, the lowering of the continuum will not allow \( n \) to increase indefinitely. For example, the magnitude of the CL, \( \Delta \epsilon \), can be estimated from the
ion sphere approximation in which the free electron density is distributed uniformly in the ion sphere. The nuclear potential is then approximately reduced by \( Z^* e/(4\pi\varepsilon_0)R_0 \) for the outer shells with \( r_n \sim R_0 \). In the cold material at solid density, \( Z^* \sim 1 \) so that \( \Delta \varepsilon \sim 10 \) eV, whereas the binding energy of an electron in \( n = 3 \) is approximately 5 eV according to the isolated AA model. Therefore, it is questionable whether \( n = 3 \) would be bound even if orbital overlapping and PI did not occur. At higher temperature, \( Z^* \) increases thereby raising \( \Delta \varepsilon \), although as the plasma expands, \( R_0 \) increases with \( \rho^{-1/3} \). Therefore, while PI may no longer be so important for higher \( n \) at the lower densities, continuum lowering may still prevent these states from being bound.

To obtain a more quantitative approximation to the number of principal quantum levels that should be used, an atomic model (Rose 1988) has been employed which solves the Schrödinger equation in a Thomas-Fermi (TF) potential. Numerical wavefunctions were calculated for orbitals of carbon atoms/ions in a range of densities and temperatures. Table 4.7.1 summarises the number of completely bound principal levels that are found for each temperature and density. The fraction of the outermost incomplete principal shell is also given because it was often found that the higher \( l \) states were not bound while the lower ones were. Thus \( 3 + 2/4 \) means that wavefunctions were found for all orbitals up to and including \( n = 3 \), but only for 2 of the 4 \( l \) states of \( n = 4 \).
Table 4.7.1. Number of bound quantum levels predicted by a Thomas-Fermi model for a carbon plasma at different temperatures and densities. $3 + 2/4$ means that wavefunctions were found for all orbitals up to and including those with principal number $n = 3$, but for only 2 of the $l$ states of $n = 4$. 

<table>
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<th>DENSITY (g cm$^{-3}$)</th>
<th>TEMPERATURE (eV)</th>
<th>1</th>
<th>10</th>
<th>30</th>
<th>50</th>
<th>75</th>
<th>100</th>
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<td></td>
<td>2 + 1/3</td>
<td>3 + 1/4</td>
<td>4 + 0</td>
<td>4 + 1/5</td>
<td>4 + 1/5</td>
<td>4 + 1/5</td>
</tr>
<tr>
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<td></td>
<td>2 + 0</td>
<td>2 + 1/3</td>
<td>3 + 0</td>
<td>3 + 0</td>
<td>3 + 1/4</td>
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<td>2 + 1/3</td>
<td>3 + 0</td>
<td>3 + 0</td>
<td>3 + 1/4</td>
</tr>
<tr>
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<td></td>
<td>1 + 1/2</td>
<td>2 + 0</td>
<td>2 + 1/3</td>
<td>2 + 1/3</td>
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</tr>
<tr>
<td>1.00</td>
<td></td>
<td>1 + 1/2</td>
<td>1 + 1/2</td>
<td>2 + 0</td>
<td>2 + 0</td>
<td>2 + 1/3</td>
<td>2 + 1/3</td>
</tr>
<tr>
<td>2.00</td>
<td></td>
<td>1 + 1/2</td>
<td>1 + 1/2</td>
<td>1 + 1/2</td>
<td>1 + 1/2</td>
<td>1 + 1/2</td>
<td>1 + 1/2</td>
</tr>
</tbody>
</table>
The results predict 2 - 3 levels to be bound as expected. Therefore, several simulations have been carried out for a 5 \( \mu \)m foil using three different values of \( N \), fixed throughout the target (\( N = 2, 3 \) and 4) and the radiation pulse described in 4.7.2 with \( T_{\text{RMAX}} = 115 \) eV. The reduction in the degeneracy of each level due to PI was not estimated in the calculation, but using different \( N \) automatically gives an assessment of the effects of extreme PI and limiting of \( N \) due to CL. It was assumed that \( \Delta \epsilon = 0 \) because it is expected to be small and its calculation is in any case uncertain.

Figure 4.7.4(a) and (b) show a comparison between \( N = 2, N = 3, (2,3) \), and \( N = 3, N = 4, (3,4) \) at two different times, 500 ps and 800 ps, (~240 ps before and ~60 ps after the peak of the radiation pulse). Note that all the solutions are very similar. The differences are largest in the cooler material but are less than 15% and 10% for (2,3) and (3,4) respectively at 500 ps. This is reduced at later times. In the hot part of the foil, the discrepancy is less than 1% for both.

In figure 4.7.5, density profiles predicted with \( N = 3 \) at 300, 500, 800 and 1000 ps are shown. Note that, as expected, the density varies from just above solid to a few times \( 10^{-2} \) gcm\(^{-3} \). Densities above solid are seen only during the first few 100 ps of the simulation. At later times, the density is approximately constant throughout the target.
Figure 4.7.4(a). Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{MAX}} = 115$ eV using 2 (----) and 3 (—) principal quantum levels. $Dm = 2 \times 10^{-5}$ gcm$^{-2}$. 
Figure 4.7.4(b). Same as 4.7.4(a) but using 3 (—) and 4 (----) principal quantum levels.
Figure 4.7.5. Predicted density profiles through a 5 \( \mu \text{m} \) N-parylene foil calculated with \( T_{\text{RMAX}} = 115 \text{ eV} \) and three principal quantum levels at 300, 500, 800 and 1000 ps. \( Dm = 2 \times 10^{-5} \text{ gcm}^{-2} \).
It is clear that under these circumstances, PI and CL have little effect on the radiation transport (provided \( n = 2 \) is not significantly affected). As expected, the effect is most pronounced in the cooler material. This is mainly because, for example, \( n = 3 \) and \( n = 4 \) have most effect on the opacity and on the populations of \( n = 1 \) and \( n = 2 \) when they are most highly populated, which occurs at low temperature. Figures 4.7.3(a), (b) and (c) clearly show the suppression of ionization due to the presence of high \( n \) at low temperature. The \( Z^* \) calculated using each \( N \) is less than that calculated by the TF model (also shown) and for any particular conditions, \( Z^* \) decreases with increasing \( N \). The values become closer as the temperature increases and density decreases, as expected. However, it is also expected that \( Z^* \) calculated by the TF model is overestimated, because with this approach the bound electron density is overestimated at the centre of the atom. This results in levels being less tightly bound and more sparsely populated and is particularly severe for the inner shells.

At higher temperature, the outer levels become less populated and the \( Z^* \) predicted by each of the \( N \) values become more nearly equal to each other (and to that predicted by the TF model) as expected. Under these conditions, the presence of say \( n = 3 \) and \( n = 4 \) has even less effect on the populations of \( n = 1 \) and \( n = 2 \). Also, the opacity below the K-edge at high temperature is dominated by f-f transitions because \( Z^* \) is high and by photoionization from \( n = 2 \) because the b-f opacity varies as \( P_n C_n^4 / n^5 \). In fact, the f-f opacity is dominant in the energy region up to \( \sim \epsilon_2 (>> \epsilon_{3,4}) \) in the hot foil. Therefore, the radiation transport solutions for the different \( N \) in the hotter, less dense plasma tend to converge.

The actual value of \( N \) will vary throughout the plasma. However,
because the choice of $N$, within reason, has little effect on the radiation transport and $Z^*$ calculated by the model is used only for the material opacity and no other plasma quantity, a constant value of $N$ is used, $N = 3$, with $\Delta \varepsilon_n = 0$. This choice is somewhat arbitrary and some experimental investigation in this area would be desirable. Using a constant $N$ has the advantage that high lying levels cannot continually oscillate between a bound and free state during the iterative solution for the $P_n$ (4.5.4) which can lead to convergence problems.
CHAPTER 5
INTERPRETATION AND INVESTIGATION OF EXPERIMENTAL DATA USING NUMERICAL MODELLING

In this chapter, the data presented in chapter 3 for a 5 μm plastic foil is analysed in detail and interpreted with the aid of computer simulations using the numerical model described in chapter 4. A global comparison of the experimental data and simulations is made. Then the measured spectral profile close to the peak of the emission is compared to the calculated profiles using three different assumed X-ray fluxes into the plastic. The effect of different L-shell opacity scaling described in 4.6 on the radiation transport and predicted spectra is then investigated. The effect of the temporal and spectral shape of the assumed emission from the gold plasma on the radiation transport and predicted spectra is assessed. Finally, before the conclusions are drawn, an investigation of non-LTE (NLTE) effects on the atomic level populations is carried out using a zero dimension (0D), time dependent, screened hydrogenic, average atom model written by Dr. S. J. Rose. To begin, a brief summary of the model is given followed by an explanation of the target heating.

Unless otherwise stated, all simulations assume an X-ray conversion efficiency towards the rear of 12% ($T_{RMAX} = 115$ eV), the $1/(hv)^2$ scaling law for the outer shells opacity, three principal quantum numbers and are for a 5 μm plastic foil. In general, the results of the simulations are presented at 500 ps and 800 ps, (240 ps and 60 ps before and after the peak of the pulse respectively).
Simulations have been carried out using a radiation transport model coupled to the 1-D Lagrangian hydrodynamics code MEDUSA (Christiansen, Ashby and Roberts 1974). The radiation transport calculation uses 116 energy groups in the 0 to 100 keV energy range with a resolution of 10 eV below 800 eV. Because the radiation field is not isotropic, the model distinguishes between inwardly and outwardly moving photons and each group is transported separately. Group averaged Planck mean opacities are calculated at the material temperature, in-line with the hydrodynamics using an average atom screened hydrogenic approximation in LTE, based on the model XSNQ (Lokke and Grasberger 1977). Only bound-free and free-free transitions are considered in the model. The Planck weighted opacity is very similar to the corresponding Rosseland value and is computationally preferable because it can be calculated analytically (4.6). Most of the calculations were performed with three principal quantum levels because numerical solutions of the Schrodinger equation in a Thomas-Fermi potential predict just three bound principal states for conditions in the foil during most of the simulation (4.7.3). Simulations with two and four principal quantum levels were carried out and showed no noticeable differences.

The behaviour of the laser irradiated thin gold layer was not calculated in the code because of the marked non-LTE behaviour of the X-ray emitting region of the gold plasma, which could not be accurately calculated by our LTE model. Therefore, several assumptions (4.7.2) are made about the emission from the gold at time t and frequency ν. Firstly, because the spectrograph is not calibrated, it is assumed that
the spectrum at the peak of the radiation pulse is Planckian, characterised by a radiation temperature $T_{\text{RMAX}}$ with additional components due to M, N and O–shell emission bands consistent with measurements (4.7.1). Secondly, it is assumed that the radiation is isotropic in the half plane because the focal spot diameter is very much larger than the target thickness. The radiative flux $F(\nu,t)$ incident on the plastic foil from the heated gold layer is then taken to be

$$F(\nu,t) = \pi B(T_{\text{RMAX}},\nu)f(\nu,t)$$

where $f$ is the measured temporal behaviour of the radiation at frequency $\nu$ and time $t$ and $B$ is the Planck function.

$T_{\text{RMAX}}$ was estimated to be 115 eV from absolute measurements which were taken under similar experimental conditions (4.7.1). This corresponds to an overall laser to X-ray conversion efficiency of approximately 12% towards the rear which is consistent with other measurements. Simulations were also carried out with values of $T_{\text{RMAX}}$ of 100 and 130 eV corresponding to X-ray conversion efficiencies of 7% and 20%, respectively.

MEDUSA assumes the ions to have a perfect gas equation of state (EOS) while the electrons obey a Thomas-Fermi EOS with modified quantum corrections to give correct solid density. Thermal electron conduction is modelled using the approximation of flux limited thermal transport. Simulations were carried out varying the transport flux limiter between 0.003 and 0.1 (5.2.3). However, no difference for the different flux limiters was observed. Hot electron heating was assumed to be negligible due to the relatively low irradiance with green laser light (4.7.3).
5.2 MATERIAL HEATING

When the X-rays are incident on the plastic foil from the gold, they are absorbed preferentially in energy regions where the material opacity is high. This occurs particularly just above photoionization edges. In this case, with plastic targets, and radiation temperatures of around 100 eV, the most important absorption occurs as a result of photoionization from the K and L-shells of the carbon atoms. Because the radiative energy tends to be deposited over approximately a mfp, the photons just above the edges are responsible for short range heating while those just below and well above the edges tend to deposit their energy over a much longer range because of the substantially increased mfp. For example, in the cold plastic target material, the mfp just below the carbon K-edge (≈ 45 Å or 280 eV) is approximately 5 μm, whereas that just above the edge is approximately 0.2 μm. Therefore, for the 5 μm plastic targets described in chapter 3, it is expected that the rear of the target will be heated, initially, predominantly by X-rays just below the K-edge. Long range heating will also occur from X-rays well above the K-edge, but significantly less radiative energy from the gold is expected to exist in this energy region. X-rays responsible for longer range heating do not heat the target as much as the short range X-rays because the energy is spread over a larger material volume. This is emphasised because the absorption varies exponentially with distance.

5.2.1 Edge Shifting

As the target is heated, the outer shell electrons of the carbon atoms are stripped rapidly and at a lower temperature than those in the K-shell. This has two effects. Firstly, the photoionization opacity due
to the outer shells decreases rapidly as the levels are depleted. This increases the photon mfp so that radiation in this energy region plays a decreasingly important role in the material heating. Secondly, the photoionization edges, in particular the K-edge, shift towards higher energy because the fewer remaining electrons become more tightly bound as electron screening decreases. This allows radiation below the shifted K-edge to propagate more deeply into the target, thereby penetrating the cooler regions where it is once more strongly absorbed. This mechanism is responsible for considerably enhanced target heating because the shifts can be very large, approximately 200 eV for cold to hydrogenic carbon. As a result, the heat wave generated by the absorbed X-rays propagates much more rapidly into the target. In figure 5.2.1, the temperature profiles in a 5 μm plastic foil at two different simulation times, 500 ps and 800 ps, are shown for a (normal) simulation in which the edges are allowed to shift and one where all the edges are held in their cold positions. As can be seen, the effect of edge shifting is cumulative and by 800 ps, the enhanced heating due to edge shifting is quite substantial (at least a factor of two) in the cooler part of the target. Also shown are the results for a simulation where just the K-edge is held in its cold position. It is easily seen that the heat wave is enhanced most by the shifting of the K-edge. On the other hand, in the cooler material, the heating is enhanced mostly by the shifting of the outer shells. This reduces the opacity for the longer range photons in the hot material allowing them to deposit more energy towards the rear of the foil.
Figure 5.2.1. Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{MAX}} = 115$ eV at 500 ps and 800 ps for different photoionization edge conditions: (---), all edges held at cold positions; (••••), K-edge only held in cold position; (----), all edges free to move. $Dm = 2 \times 10^{-5}$ gcm$^{-2}$. 
5.2.2 Edge Smearing

Like edge shifting, edge smearing results from heating of the target material. However, edge smearing occurs because the target is heated unevenly by the X-rays so that temperature (and density) gradients are set up in the foil. Along the temperature gradient, different ionization stages will predominate whose edge positions are situated at different energies. The result is that the opacity spectrum in the region of the edges is smeared out over the temperature gradient and becomes much flatter. This is demonstrated in figure 5.2.2 where the total attenuation due to three 1 µm slabs of mylar (C_{10}O_{4}H_{8}), at different temperatures, is shown for conditions similar to those experienced in the heated foils, 1, 50 and 100 eV respectively and 0.1 g cm\(^{-3}\) in the 0.2 – 1.0 keV energy range. Also shown is the attenuation due to a cold slab with the same total mass as the three 1 µm samples. The opacity spectrum has been smeared out but is always below the cold spectrum. The result is that the overall attenuation of radiation at each photon energy traversing the temperature gradient is more nearly equal in this energy range. In turn, this affects the way in which the material is heated, which can be seen by comparing the shapes of the radiative energy deposition rate profiles shown in figure 5.2.3 for simulations with and without edge shifting and therefore smearing. It can be clearly seen that the two profiles are distinctly different. When edge shifting and thus smearing is allowed, the heating profile is much flatter along the temperature gradient (see figure 5.2.1) as expected for a smeared out opacity spectrum. The deeper penetration of the profile calculated with edge shifting is due to shifting and not smearing.
Figure 5.2.2. Total attenuation due to three 1 μm slabs of mylar (C₁₀O₄H₈) at a density of 0.1 gcm⁻³ and temperatures of 1, 50 and 100 eV in the 0.3 - 0.6 keV energy region (---). A, B and C are due to the carbon K-edge at 1, 50 and 100 eV respectively. D is due to the cold oxygen K-edge. The oxygen K-edge produced at higher temperatures lies outside the range of the graph. Also shown is the attenuation due to a single slab of the cold material (-----) containing the same mass as the three 1 μm samples.
Figure 5.2.3. Power deposition profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV calculated with fixed photoionization edge positions and no smearing (—) and free photoionization edges and smearing (----). $Dm = 2 \times 10^{-5} \text{ gcm}^{-2}$. 
In the model, both edge shifting and smearing are automatically taken into account. This is not the case when opacity data is tabulated prior to a simulation and care should be exercised when using tabulated data. Tabulation of opacity data, although often necessary, also has the unfortunate tendency of forcing edges to exist at the energies calculated at the tabulated data points straddling the actual plasma conditions.

5.2.3 Electron Thermal Transport vs Radiation Transport
Several simulations were carried out varying the classical electron flux limiter, $f_e$, between 0.003 and 0.1. Identical results were obtained irrespective of the choice of $f_e$. This implies that the temperature gradients set up in the foil are not sufficient to cause the Spitzer heat flow to exceed 0.003 times the free streaming limit. A further simulation was performed permitting no electron thermal conduction to occur between cells. Very similar results were also obtained. Figure 5.2.4 shows the foil temperature profiles at $t = 300$ ps and $t = 500$ ps for a simulation with and without electron thermal conduction. The difference in the results are at most 4% but generally around 1% at $t = 500$ ps. At the later time, the discrepancy is less than 1%. This implies that the major energy transport mechanism in the foils is radiation transport. This need not necessarily be the case if there were a strong thermal electron source next to the plastic which is not taken into account in the model. This is not expected to be the case in the experiment because it is estimated that the ratio of the radiation to thermal conductivities in the gold is around $10^4$ (4.7.1).
Figure 5.2.4. Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV at 300 ps and 500 ps with electron thermal conduction (-----) and without electron thermal conduction (—). $Dm = 2 \times 10^{-5}$ gcm$^{-2}$. 
5.2.4 Importance of Radiation Transport

By radiation transport, it is meant that photons are absorbed and re-emitted at least once. To assess the importance of this process on the material heating, a further standard simulation was carried out which allowed the re-emission of photons but no further absorption, i.e. the re-emitted radiation was allowed to escape from the system. This is similar to an approximation which is often made in laser plasma hydrocodes which do not include radiation transport calculations but which allow the emission due to thermal bremsstrahlung to escape from the plasma. In figure 5.2.5, the temperature profiles at simulation times 500 ps and 800 ps are compared for calculations with and without absorption of emitted radiation. Clearly, the re-emission of radiation plays a very important role in the material heating since significant cooling is observed when the re-emitted photons are allowed to escape. The absorption of these photons is important because it is this which is at least partially responsible for heating the target to give the normal profiles; partially responsible because the enhanced heating also allows deeper penetration of the radiation incident on the plastic. The heating rates in the target are compared for the simulation at 500 ps in figure 5.2.6 where the power deposition profiles are shown. It is seen that the heating penetrates more deeply into the target when the radiation transport is included. Also, the energy absorption rate in the hotter part of the plasma is increased by up to approximately 50%. Like edge shifting, the enhanced heating due to the radiation transport is cumulative because this allows deeper penetration of the heating radiation, which in turn encourages the radiation transport. For example, at 500 ps, the difference in the temperature profiles is at most 30 – 40% whereas, at 800 ps, differences in excess of 50% are observed.
Figure 5.2.5. Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV at 500 ps and 800 ps with absorption of reemitted radiation (---) and allowing reemitted radiation to escape (—). $Dm = 2 \times 10^{-5}$ g cm$^{-2}$. 

Lagrangian Coordinate (Dm)

Electron Temperature (eV)

- 800 ps
- 500 ps

Figure 5.2.5. Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV at 500 ps and 800 ps with absorption of reemitted radiation (---) and allowing reemitted radiation to escape (—). $Dm = 2 \times 10^{-5}$ g cm$^{-2}$. 

Lagrangian Coordinate (Dm)

Electron Temperature (eV)

- 800 ps
- 500 ps

Figure 5.2.5. Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV at 500 ps and 800 ps with absorption of reemitted radiation (---) and allowing reemitted radiation to escape (—). $Dm = 2 \times 10^{-5}$ g cm$^{-2}$. 

Lagrangian Coordinate (Dm)

Electron Temperature (eV)

- 800 ps
- 500 ps

Figure 5.2.5. Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV at 500 ps and 800 ps with absorption of reemitted radiation (---) and allowing reemitted radiation to escape (—). $Dm = 2 \times 10^{-5}$ g cm$^{-2}$. 

Lagrangian Coordinate (Dm)

Electron Temperature (eV)

- 800 ps
- 500 ps
Figure 5.2.6. Predicted power deposition profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV with absorption of reemitted radiation (-----) and allowing reemitted radiation to escape (—). $D_m = 2 \times 10^{-5}$ gcm$^{-2}$. 
5.3 COMPARISON OF PREDICTED SPECTRA WITH EXPERIMENTAL DATA FOR A 5 \( \mu \)m PLASTIC FOIL

5.3.1 Global Comparison
Figure 5.3.1(b) shows a 3-D plot of the predicted intensity of radiation in units of Wm\(^{-2}\)A\(^{-1}\) into 2\( \pi \) steradians transported through a 5 \( \mu \)m plastic foil target as a function of time and wavelength for \( T_{\text{RMAX}} = 115 \) eV. At each time spectral emission profiles as a function of wavelength were obtained from a 1-D integration of the radiative transfer equation through the heated thin foil target at the mid point of each energy group in the 10 to 70 A wavelength region. This can be seen more clearly in figure 5.3.1(c) where the vertical scale has been expanded. In common with the experiment, the simulations show a clear edge which is due to the carbon K-edge, the early turn on of radiation just below and well above the edge, progressively retarded emission of radiation towards longer wavelength away from the edge and a gradual shift in the edge position as the foil is heated. Figure 5.3.1(a) shows the soft X-ray spectrum assumed to be emitted from the gold layer before it has passed through the plastic foil corresponding to our best estimate of the experimental conditions for the data shown in figure 3.4.5(b) (\( T_{\text{RMAX}} = 115 \) eV).

5.3.2 Spectral Comparison At Peak Emission
Figure 5.3.2 shows the variation of the intensity as a function of wavelength of the experimental data of figure 3.4.5(b) close to the peak emission. The shaded band shown results from the envelope formed by several densitometer traces taken within \( \pm 150 \) ps of the peak of the emission.
Figure 5.3.1(a). 3-D plot of the spectrum incident on the 5 μm N-parylene foil as a function of time and wavelength. The spectral distribution is Planckian at the peak of the pulse with a temperature of 115 eV. The vertical scale is linear and has units of Wm$^{-2}$Å$^{-1}$ into 2π str with the base of the graph drawn at one hundredth the maximum intensity plotted.
Figure 5.3.1(b). 3-D plot of the predicted soft X-ray spectrum emitted from the rear of a 5 \( \mu \text{m} \) N-parylene foil for the incident spectrum in 5.3.1(a). The vertical scale is the same as in 5.3.1(a).
Figure 5.3.1(c). Same as 5.3.1(b) but the vertical scale is expanded. The base is drawn at one twentieth the maximum intensity plotted.
Figure 5.3.2. Experimentally measured soft X-ray intensity versus wavelength close to the peak of the radiation pulse ($\pm 150$ ps) from the rear of a $5\mu m$ N-parylene foil irradiated by a $1.5 \times 10^{14} \text{ Wcm}^{-2}$, 700 ps (FWHM) green laser pulse. The band results from the variation of the signal during the considered 300 ps interval and experimental uncertainties.
The time interval of $\pm 150$ ps is the estimated uncertainty in the peak of the emission from the gold as the time resolved spectrograph was not absolutely timed. The relative response of the instrument was estimated by taking the various components such as filters, mirror, grating and streak camera photocathode into consideration.

In comparison with the experimental data of figure 5.3.2, the predicted rear side spectra from the rear of a 5 $\mu$m plastic foil target at the peak of the radiation pulse for three different values of $T_{RMAX}$ (100, 115 and 130 eV) are shown in figure 5.3.3. Note the significant shift in the carbon K-edge for the different temperatures. In addition, the assumed soft X-ray spectrum for 115 eV incident on the plastic target is shown in figure 5.3.3.

The position of the edge feature in the experimental data is at approximately 45 A which corresponds to the cold carbon K-edge position. Maximum shifts of about 1 A were seen during the heating phase. In contrast, much larger shifts of the carbon K-edge positions are predicted in the simulations depending on the temperature of the heated material. For example, a shift of more than 10 A towards shorter wavelength is seen for a peak radiation temperature of 115 eV.

Although the code calculations are in overall agreement with the experiment, the small apparent shift in the experimental edge position is not predicted by any of the values of $T_{RMAX}$ in figure 5.3.3. The small shift is however predicted if a value of $T_{RMAX} = 75$ eV is used. In this case the radiation does not burn through a 5 $\mu$m thick target, leaving a layer of cold material towards the rear of the foil. The cold layer attenuates strongly in the wavelength region between 24 and 44 A which is contradictory to the experimental observations.
Figure 5.3.3. Predicted soft X-ray spectrum from the rear of a 5 µm N-parylene foil for values of TRMAX of 100 (---), 115 (---) (curve a) and 130 eV (----) close to the peak of the emission. The curve labelled (b) is the incident spectrum for a temperature of 115 eV. The vertical scale is in absolute units Wm⁻²Å⁻¹ into 2π str.
In addition, a radiation temperature of 75 eV corresponds to a conversion efficiency of only 2% which is well below the conversion efficiencies published for green laser light at an irradiance of around $10^{14}$ W cm$^{-2}$.

It is believed that $T_{\text{RMAX}} = 115$ eV is most realistic and suggested that the small shift in the edge position is due to the filling in of the photoabsorption below the shifted K-edge by line absorption. The edge in figure 5.3.2 would then correspond to the side of a strong carbon absorption line. Figure 5.3.4 shows the degree of ionization through a 5 μm foil predicted by the model at 800 ps for $T_{\text{RMAX}} = 115$ eV. The shaded band shows the spread of the results calculated using two, three and four principal quantum levels. It can be clearly seen that there are significant fractions of H - and He - like ions throughout the entire foil. Towards the rear of the target, significant fractions of even Li - like ions can be expected to exist. Figure 5.3.5 shows a trace across wavelength for the data presented in figure 3.4.5(b) for a 5 μm foil close to the peak emission. Also shown are the positions of the first three resonance lines of H - and He - like carbon ions (wavelength data taken from Kelly 1987). The accuracy of the trace with respect to the wavelength scale is estimated to be ± 1 Å. Although positive identification of absorption or emission features due to these lines is not conclusive, it is clear that, in particular, the He - like transitions are very important in determining the spectrum just above (in energy) the edge feature.
Figure 5.3.4. Predicted degree of ionization through a 5 \( \mu \text{m} \) N-parlylene foil at 800 ps for \( T_{\text{RMAX}} = 115 \) eV. The shaded band shows the spread of results from calculations using 2, 3 and 4 principal quantum levels.
Figure 5.3.5. Densitometer trace across the spectrum close to the peak emission from the rear of a gold coated 5 μm N-parylene foil irradiated by a $1.5 \times 10^{14}$ W cm$^{-2}$, 700 ps (FWHM) green laser pulse. The positions of the first three resonance transitions of $^{\text{He}}$- and $^{\text{H}}$-like carbon ions are shown.*

* taken from Kelly (1987).
Improvements to the opacity calculations to include both the distribution of ionic states around the average in both the bound-free and bound-bound transitions is underway to assess whether this is the correct explanation for the discrepancy. A further possible explanation for the discrepancy is that the bright feature below 45 A in figure 5.3.2 originates from the gold plasma. This would correspond to emission from the gold O-shell (Eidmann and Kishimoto 1986). However, time resolved spectra of gold emission through ultra-thin plastic foil targets 0.1 μm thick, which are too thin to significantly affect the transmitted radiation, suggest that this is not the case (3.4).

5.4 \(1/(hv)^2\) vs \(1/(hv)^3\) L-SHELL OPACITY SCALING

It was seen in 4.6 that in the cold material, the carbon L-shell opacity scaled more nearly as \(1/(hv)^2\) rather than \(1/(hv)^3\) which was attributed to the non-Coulombic nature of the potential in the outer regions of the atom. When the standard cubic scaling is used the cold L-shell opacity is calculated to be \(1.04 \times 10^2 \text{ cm}^2\text{g}^{-1}\) just below the carbon K-edge, as opposed to \(1.74 \times 10^3 \text{ cm}^2\text{g}^{-1}\) when the modified square scaling is used. This results in an increase in the photon mfp by more than an order of magnitude for the \(1/(hv)^3\) law in this region from around 5 μm to around 80 μm. This will have virtually no effect on the K-shell opacity because it is very much larger than that due to the L-shell (above the K-edge). Therefore, it is expected that the predicted spectral results just below the K-edge will differ significantly when the different scalings are used, at least until the foil becomes hot and the outer electrons are stripped.
5.4.1 Effect on Radiation Transport

On the other hand, it is expected that the difference in the material heating will not be so significant. This is because, although the increased opacity below the K-edge due to the $1/(hv)^2$ scaling will increase the material heating, the electrons in the outer shells are rapidly stripped as the plasma temperature rises. Also, following the discussion of 5.2, the radiation in this region heats the matter over a larger volume than for example radiation just above the K-edge and in comparison, contributes less to the overall heating of the material. More heating is obtained from radiation closer to the L-edge but there the opacities predicted by the two scaling laws are more nearly equal anyway. In addition, it was found in 5.2 that absorption of radiation above the cold K-edge was responsible for a large fraction of the material heating. Figure 5.4.1 shows temperature profiles for the two different scalings at two different simulation times, 500 ps and 800 ps. The enhanced heating, particularly in the cooler material, is clearly observed for the $1/(hv)^2$ law. The difference is greatest (~ 30%) in the cooler but not cold (only ~ 10%) plasma. The increased heating due to the $1/(hv)^2$ outer shells opacity scaling law is also responsible for increased K-edge shifting, particularly in the higher temperature plasma. This allows deeper penetration of radiation just above the cold K-edge position into the cooler plasma where it is strongly absorbed contributing to the enhanced heating seen there. However, radiation in this energy region heats over a shorter range not significantly affecting the cold plasma much deeper in the target (5.2).

This is also reflected in the spectrum close to the peak of the pulse at 800 ps shown in figure 5.4.2. The spectra are very similar here because, by this time, the foil is significantly heated throughout and the
Figure 5.4.1. Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV at 500 ps and 800 ps using standard $1/(\hbar \nu)^3$ outershells opacity scaling (---) and the modified $1/(\hbar \nu)^2$ law (----). $Dm = 2 \times 10^{-5}$ g cm$^{-2}$. 
Figure 5.4.2. Predicted soft X-ray spectrum from the rear of a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV with standard $1/(h\nu)^3$ outer shells opacity scaling (—) and the modified $1/(h\nu)^2$ law (----).
occupancy of the outer atomic levels is small. In fact, if L is the
target thickness at 800 ps, the mfp just below the shifted K-edge is 5L
and 15L for the square and cubic scaling laws respectively. Therefore,
the spectra in this energy region are very similar. The edge positions
predicted by the two laws are different but only by around 2 – 3 Å.
This is consistent with the slightly different degrees of heating.

5.4.2 Effect on Temporal Evolution of Spectrum.
In contrast, a considerable difference is observed when comparing the
temporal evolution of the predicted spectra at a specific wavelength
below the cold K-edge. Figure 5.4.3(a) shows the temporal evolution
of the predicted rear side spectra at 48 Å (255 eV) for both scaling
laws. The large difference at early times is due to the very different
mfps resulting from the different scalings. The assumed variation of
the spectral emission from the gold plasma is also shown together with
the experimentally measured rear side emission at this wavelength. The
measured profile has been normalised to the profile calculated using
the $1/(h\nu)^3$ law and it has been assumed that the peaks in the
measured and predicted profiles coincide. Therefore, although the
square law fits the cold opacity data best, it appears that the better fit
with the experimental data is obtained from the $1/(h\nu)^{3}$ law. This is
difficult to understand because at, for example 500 ps, figure 5.4.1
shows that there is still a significant amount of cold material towards
the rear of the foil. The model predicts that $Z^* = 1$ for just over 1/3
of the plasma concentrated at the rear of the target. Therefore, if
the model is correct, it should be expected that the $1/(h\nu)^2$ scaling
gives the best agreement at this time, which is not the case in figure
5.4.2. However, the calculation is expected to work least well at low
temperature and ionization. In addition, it will be seen that the
Figure 5.4.3(a). Temporal evolution of predicted emission at 48 Å (255 eV) from the rear of a 5 μm N-parylene foil for $T_{RMAX} = 115$ eV with standard $1/(h\nu)^3$ outer shells opacity scaling (—) and the modified $1/(h\nu)^2$ law (----). Also shown is the measured profile (- - -) and the assumed profile of the emission from the gold (-----).
assumed radiation pulse shape also affects the rear side temporal profiles (5.5).

As expected, above the cold K-edge the temporal profiles agree closely for both scalings. Figures 5.4.3 (b) and (c) show the same quantities plotted in 5.4.3(a) for 39 A (320 eV, just above the cold K-edge) and 15 A (830 eV) respectively. The agreement here with the experiment is less good, although at 15 A this is probably due to the assumed temporal shape of the gold spectrum (5.5.2). At 39 A the discrepancy cannot be explained by the assumed pulse shape. The late and rapid turn on of the emission here can be explained as follows. The short mfp above the edge prevents significant emission until the edge has shifted to some higher energy throughout most of the foil. Hence a rapid turn on of radiation occurs when the last but a few mfps become significantly heated by the radiation heat wave for the edge to have shifted to above 39 A (320 eV). The rapid turn on is not observed in the experiment and is discussed further in 5.5.2.

5.5 EFFECT ON MATERIAL HEATING AND PREDICTED SPECTRA OF RADIATION PULSE SHAPE

5.5.1 Spectral
To assess the individual effects of spectral features such as M, N and O emission bands on the material heating, simulations were carried out with approximate spectral forms for M, N and O–bands individually. These features were included in such a way that the total power in the radiation pulse at any given time is constant irrespective of the spectral form (see 4.7.2).
Figure 5.4.3(b). Same as 5.4.3(a) but at 39 Å (320 eV).
Figure 5.4.3(c). Same as 5.4.3(a) but at 15 Å (830 eV).
The M-band was simulated by adding 4% (Siegel 1987) of the total X-ray power to four energy groups centred around 2.8 keV and extending from 2.2 – 3.5 keV (LLE Review 1984). The enhanced emission in the N and O-bands was simulated by adding 10% of the total X-ray energy to groups extending from 500 – 800 eV and 150 – 300 eV respectively (Eidmann and Kishimoto 1986). For convenience and without better experimental evidence, the spectral form of the extra energy in these bands was taken to be

\[ f(\nu)_{\text{N, O}} = C \nu^{m-1} \exp(-\nu^m) \]

because it is easily integrable between any limits and allows the shape, extent and position of the band to be changed very easily. \( C \) is a normalisation constant which is found by integrating the above expression over the band interval. It was found that altering the spread of the energy within the boundaries of the emission bands by using different \( m \) values produced no noticeable differences in the results.

In figures 5.5.1(a), (b) and (c), temperature profiles taken at two different times, 500 ps and 800 ps, calculated with and without M-band, N-band \((m = 3)\) and O-band \((m = 3)\) emissions are shown. The results are very similar and the slight differences are easily explained. When the M-band is included, the temperature profile appears to lag behind the standard profile. This is because the mfp even in the cold material for radiation in the M-band is \( \approx 30 \mu \text{m} \). Therefore, the M-band heats the whole foil approximately evenly and instantaneously, but also very little because the mfp is much greater than the target thickness. A similar effect would be achieved simply by subtracting 4% of the energy from the pulse. The difference in the profiles is at most 8% in the cooler material.
Figure 5.5.1(a). Predicted temperature profiles through a 5 μm N-parylene foil for $T_{RMAX} = 115$ eV at 500 ps and 800 ps with 4% M-band emission included (—) and without M-band emission (----). $Dm = 2 \times 10^{-5}$ gcm$^{-2}$. 
Figure 5.5.1(b). Same as 5.5.1(a) but with 10% N-band instead of M-band emission included.
Figure 5.5.1(c). Same as 5.5.1(a) but with 10% O-band instead of M-band emission included.
The inclusion of the N-band increases the temperature in the hot material, but decreases it somewhat in the cooler regions of the target. In contrast, the O-band has the opposite effect. This results from the different positions of the extra emission features in the photon energy spectrum. The O-shell is situated just below the cold K-edge. Following the discussion of 5.2 and 5.3, the radiation in this region is responsible for long range heating. Therefore, when the O-band is included, enhanced heating (~ 10%) is observed in the cooler plasma. In the hot foil, there is a slight reduction in the material temperature (~ 2%). This is because the energy in the O-band (long range heating region) has been transferred from the rest of the spectrum, some of which will have come from the short range heating regions.

In contrast, the N-band is situated above the cold and indeed hot (~ 480 eV) K-edge position. As a result, the hot plasma is heated more (~ 2%) because some energy has been transferred to a short range heating region from the longer range heating regions. On the other hand, in the cool plasma, the heating is less (~ 2 - 3%) both because there is less energy in the longer range heating regions and because this energy has been transferred to the N-band where it is more strongly absorbed in the hot plasma.

The difference in the spectral emissions is minor and local, e.g. enhanced O-shell emission in the gold spectrum is observed as enhanced emission in the same energy region in the predicted rear side spectrum with little further differences. This is demonstrated in figures 5.5.2(a), (b) and (c), (d), where the assumed gold spectra with N and O emission bands and the predicted rear side spectra are shown as a function of time and wavelength.
Figure 5.5.2(a). 3-D plot of the spectrum incident on the 5 μm N-parylene foil as a function of time and wavelength. The spike is due to the assumed gold N-band emission. The integrated spectral flux at peak emission is described by a temperature of 115 eV. The vertical scale is linear and has units Wm^{-2}A^{-1} into 2π str with the base of the graph drawn at one hundredth the maximum intensity plotted.
Figure 5.5.2(b). 3-D plot of the predicted soft X-ray spectrum emitted from the rear of a 5 μm N-parylene foil for the incident spectrum in 5.5.2(a). The vertical scale is the same as in 5.5.2(a).
Figure 5.5.2(c). Same as for 5.5.2(a) but with gold O-band instead of N-band emission included.
Figure 5.5.2(d). Same as for 5.5.2(b) but for incident spectrum shown in 5.5.2(c).
5.5.2 Temporal

To assess the effect of the temporal form of the assumed radiation pulse on the radiation transport and predicted spectra, the rising part of the pulse was changed from a Gaussian to a supergaussian. However, the FWHM remained 700 ps and the total energy delivered by the pulse was kept the same. The effect of this is to deliver the energy at a higher rate during the pulse rise. The change in the pulse shape shifted the peak from 740 ps to 585 ps. This is shown in figure 5.5.3.

It was found that the results for the two pulse shapes were very similar provided the comparison was made when the same amounts of energy had been delivered and not when the instantaneous radiative powers were equal. Figure 5.5.4(a) shows temperature profiles taken at 500 ps for the new pulse and 600 and 700 ps for the old pulse. The energy supplied by the new pulse in this time (6.04 \(10^7\) Jm\(^{-2}\)) is approximately half way between that supplied by the old pulse at 600 and 700 ps (5.24 \(10^7\) Jm\(^{-2}\), 6.98 \(10^7\) Jm\(^{-2}\)). The temperature profile calculated by the new pulse at 500 ps is of exactly the same form as those for the old pulse at 600 and 700 ps and very nearly half way between the two. Also shown are profiles from the new and old pulses at 300 ps and 400 ps respectively. The energies supplied by the new and old pulses in these times are 2.57 \(10^7\) Jm\(^{-2}\) and 2.33 \(10^7\) Jm\(^{-2}\) respectively (see 5.5.4(b)). Again, it is clear that the profiles are of the same form, with the old pulse profile lagging the new pulse profile because of the smaller supplied energy.
Figure 5.5.3. Comparison between the analytic fit (----) to the experimentally measured temporal profile of the gold emission and a faster rising pulse with the same area (—).
Figure 5.5.4(a). Predicted temperature profiles through a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV using standard temporal profile (----) at 400 ps, 600 ps and 700 ps. Also shown are profiles predicted using the faster rising pulse at 300 ps and 500 ps (---). $Dm = 2 \times 10^{-5}$ gcm$^{-2}$
Figure 5.5.4(b). Energy supplied versus time for the standard (-----) and the faster rising (----) pulses compared in figure 5.5.3.
These results imply that the material heating in this case is primarily a function of the supplied energy and not the rate at which it is delivered. The reason for this is as follows. The rate of cooling that results from hydrodynamic expansion or the escape of radiative energy from the foil is small compared to the rate at which the foil is heated. Figure 5.5.5 shows the power in the old radiation pulse and the powers radiated from the front and rear of the target as a function of time. Clearly, during the heating phase, the rate of cooling due to the escape of radiation from the plasma is small compared to the heating rate. This is because thermodynamically, during the pulse rise where the radiation temperature of the spectrum from the gold is either rising or constant, it is impossible for the radiation temperature in any part of the foil to exceed the temperature of the heating radiation. Consequently, the material temperature is always less than that of the radiation pulse. Recalling equation 2.2.2, the monochromatic emission from plasma with optical depth $\Delta \tau_\nu$ in LTE is given by

$$B_\nu [1 - \exp(-\Delta \tau_\nu)] \leq B_\nu$$

Therefore, because $B_\nu(T_1) > B_\nu(T_2)$ for $T_1 > T_2$ for all $\nu$, the integrated emission from the foil is always less than that injected by the radiation pulse. In addition, for the equality to be approximately valid in the above expression, the optical depth must be very large. However, high temperature and optical depth are not in general compatible in the foil. Figure 5.5.5 shows that before burn through at around 600 ps, the total power emitted from the foil is approximately an order of magnitude less than that supplied. Once the energy has been delivered to the foil, it is transported instantaneously compared to the hydrodynamic time scale and retarded effects are negligible. Also, radiation transport has already been seen to be the major energy
Figure 5.5.5. Predicted temporal profiles of the total radiative power emitted from the front $P_{of}$ and the rear $P_{ob}$ of a 5 μm N-parylene foil for $T_{\text{RMAX}} = 115$ eV. The incident heating power profile $P_{i}$ is also shown.
transport mechanism in the plasma (5.2.3 and 4). Therefore, very similar transport and heating solutions will result from different pulse shapes irrespective of the rate at which the energy is delivered, provided the amount supplied is the same. This is only true if the time scales are not significantly different. For example, if the energy were delivered in a time much less than the hydrodynamics time scale for one pulse and on a hydrodynamics time scale for the other, then different results would be expected.

In contrast, while the pulse shape has little effect on the radiation transport and heating, provided the energy supplied is considered and not instantaneous radiative power, the temporal shape of the predicted rear side spectrum is affected by the temporal profile of the incident spectrum as expected. In figures 5.5.6(a),(b) and (c), the temporal profiles of the rear side emission from the foil are compared for the two pulses at the same wavelengths as in 5.4, 48 A (255 eV), 39 A (320 eV) and 15 A (830 eV). The input spectra for the two pulses are also shown. Quite clearly, at 48 A and 15 A, the different pulses produce distinctly different temporal profiles for the rear side spectra which reflect the shape of the incident spectrum. In contrast, little difference in the forms of the profiles is seen at 39 A. Therefore, while the temporal form of the spectrum is critically dependent on the shape of the assumed incident spectrum away from the K-edge, significantly less dependence is seen just above the edge at 39 A. This is because burn through of the foil must occur before an emitted spectrum is predicted in this energy region (5.4). By the time it does for this case, the power in the two radiation pulses is very similar.
Figure 5.5.6(a). Temporal evolution of predicted emission at 48 Å (255 eV) from the rear of a 5 μm N-parylene foil for the standard radiation pulse (----) and the faster rising pulse (—). The incident profiles are also shown.
Figure 5.5.6(b). Same as 5.5.6(a) but at 39 A (320 eV).
Figure 5.5.6(c). Same as 5.5.6(a) but at 15 A (830 eV).
In 5.4, it appeared that the standard \( 1/(h\nu)^3 \) outer shell opacity scaling agreed better with the experimental data just below the K-edge, even though it was expected that the \( 1/(h\nu)^2 \) law should produce better agreement. Here, it is seen that any discrepancy between the experimental results and the \( 1/(h\nu)^2 \) scaling would be exaggerated by deviation from the assumed profile of the radiation pulse shape (including width) from the gold due to variation in the laser pulse. It is believed that this is largely responsible for the discrepancy between experiment and theory at 15 A and at least partially responsible for that at 48 A. However, this cannot explain the significant differences observed between the predicted and measured profiles just above the K-edge.

5.6 NON-LTE EFFECTS

The validity of the assumption of LTE in the foil plasmas has been assessed using a 0D, time dependent, non-LTE (NLTE), screened hydrogenic, average atom model, written by Dr. S. J. Rose, based on XSNQ-U (Lokke and Grasberger 1977). The conditions calculated at the front, centre and rear of a 5 \( \mu \)m foil using the standard simulation and the LTE model were used in the NLTE calculation to obtain NLTE AA level populations. The radiation field at each point was assumed to be Planckian, characterised by a temperature, \( T_R \), calculated from the sum of the total fluxes incident on the cell from each side i.e.

\[
\sigma T_R^4 = \sum F \quad \text{5.6.1}
\]

The 0D model assumes that the radiation field is isotropic so that \( \sigma T_R^4 \)
passes each way through unit area into both halves of a sphere. Therefore, the temperature used to describe the radiation field is given by $2^{-1/4}$ times that defined above. Calculations were performed both with and without a radiation field present. When no radiation field was included, radiative coupling between levels was via spontaneous decay only. It was found that the foil plasmas are very well described by the LTE approximation and significantly that the presence of the radiation field away from the rear of the target was responsible for maintaining the populations closer to LTE than when no radiation field was present.

5.6.1 Effect of Radiation Field

In figures 5.6.1(a),(b),(c), 5.6.2(a),(b),(c) and 5.6.3(a),(b),(c), the NLTE populations calculated as a function of time are compared with and without the inclusion of a radiation field for conditions at the front, centre and rear of the target. Also shown for reference are the populations calculated by the LTE model during the simulation. Here, it is simply noted that, on the whole, there appears to be a very good agreement between the LTE model and NLTE calculation. Differences at early times in $n = 2$ and $n = 3$ are due to the initialisation of the calculation at low temperature when $Z^*$ is small.

The populations of $n = 1$ calculated with a radiation field in general tend to be less than those calculated without. On the other hand, the opposite appears to be true for $n = 2$ and $n = 3$. The reason for this is that the radiation field tends to photopump electrons from $n = 1$ to $n = 2$, $n = 3$ and the continuum. However, it is expected that collisional coupling between $n = 2$ and $n = 3$ and the continuum is more effective than between $n = 1$ and any other state.
Figure 5.6.1(a). Level populations versus time for $n = 1$ calculated at the front of a 5 μm N-parylene foil for $T_{RMAX} = 115$ eV using the LTE model (---), the NLTE model with a Planckian radiation field (o) and the NLTE model without a radiation field (×).
Figure 5.6.1(b). Same as 5.6.1(a) but at the centre of the foil.
Figure 5.6.1(c). Same as 5.6.1(a) but at the rear of the foil.
Figure 5.6.2(a). Same as 5.6.1(a) but for $n = 2$. 

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Figure 5.6.2(b). Same as 5.6.1(b) but for n = 2.
Figure 5.6.2(c). Same as 5.6.1(c) but for $n = 2$. 
Figure 5.6.3(a). Same as 5.6.1(a) but for \( n = 3 \).
Figure 5.6.3(b). Same as 5.6.1(b) but for $n = 3$. 
Figure 5.6.3(c). Same as 5.6.1(c) but for $n = 3$. 
Therefore, $n = 2$ and $n = 3$ are expected to depart from LTE less and the photopumping is seen to deplete $n = 1$ and slightly enhance the populations of $n = 2$ and $n = 3$.

5.6.2 Comparison of NLTE and LTE Populations

To compare the NLTE calculation with LTE predictions, LTE populations must be calculated from the NLTE energy levels i.e. from the same atomic potential. These are in general different for calculations with and without a radiation field. Recalling the discussion of 2.1.1, when a system is in LTE all collisional processes occur at their equilibrium rates. If the radiation field is Planckian with the same temperature as the material, ($T_R = T_M$), then all the radiative processes will occur at their equilibrium rates also (at $T_M$) and there will be no departure from LTE. If, however, $T_R > T_M$ or $T_R < T_M$ and collisional processes do not dominate over radiative processes, then departures from LTE are inevitable. Furthermore, even if $T_R = T_M$ and instantaneously collisional processes do dominate, there is no guarantee that the populations are described by their LTE values because the system may be time dependent.

In the foil plasmas, some departures from LTE are expected because in general $T_R \neq T_M$ and both are time dependent. Mostly, the departures from LTE are largest for $n = 1$ for the reasons discussed in 5.6.1 and 2.1.1. In figures 5.6.4(a), 5.6.5(a) and 5.6.6(a), the temporal variations of the material and radiation temperatures ($2^{-1/4} T_R$ defined in 5.6.1) are shown at the front, centre and rear of the target respectively. In figures 5.6.4(b), 5.6.5(b) and 5.6.6(b), a factor $\delta(n)$ defined by

$$\delta(n) = \frac{(P_n^{\text{NLTE}} - P_n^{\text{LTE}})}{P_n^{\text{NLTE}}}$$
(where \( P_n \) is the population of level \( n \)) is shown when no radiation field is present (top) and when the Planckian radiation field is included (bottom), for \( n = 1 \) and conditions experienced at the front, centre and rear of the target respectively. A negative \( \delta(1) \) indicates depletion of \( n = 1 \) by photopumping. In general, the signs of \( \delta(n) \) for \( n = 2 \) and \( n = 3 \) are exactly opposite to those for \( n = 1 \). Note that the vertical scales for \( \delta(1) \) vary from plot to plot and often from -ve to +ve values on the same graph.

Consider first the results for \( \delta(1) \) calculated without a radiation field at the front of the foil. There is a steady departure from LTE with \( n = 1 \) becoming overpopulated by up to 20\%. The peak in the departure from LTE occurs approximately 200 ps after the peak in \( T_R \) and \( T_M \) (\( = T_e \)). This is due both to the continual decompression of the target and the time dependence of the system. \( \delta(1) \) is reduced at later times because of the dramatic decrease in \( T_e \).

The general pattern is repeated at the centre and rear of the foil although the departures from LTE here are generally less (15\% centre, 0.5\% rear). This is largely due to the significantly lower \( T_e \) encountered. It is not a density effect because, particularly at late times (\( \gtrsim 1000 \) ps), the density throughout the foil is approximately constant. If anything, the plasma density decreases towards the rear of the foil (see figure 4.7.5). At the centre of the foil, the delay between \( \text{MAX}(\delta(1)) \) and \( \text{MAX}(T_e) \) is only 50 ps, whereas there appears to be no delay at the rear of the target. Note that at the back of the foil, \( T_e \) falls off very slowly after the peak at around 1050 ps and the density continues to fall steadily after this time which suggests that \( \delta(1) \) should increase or remain approximately constant.
Figure 5.6.4(a). Predicted electron and radiation temperatures versus time at the front of a 5μm N-parylene foil for $T_{RMAX} = 115$ eV.
Figure 5.6.4(b). Predicted percentage difference versus time between NLTE and LTE populations of $n = 1$ calculated without an ambient radiation field (top) and with a Planckian radiation field (bottom) characterised by the radiation temperature in figure 5.6.4(a), at the front of a 5μm N-parylene foil for $T_{\text{Rmax}} = 115$ eV.
Figure 5.6.5(a). Same as 5.6.4(a) but at the foil centre.
Figure 5.6.5(b). Same as 5.6.4(b) but at the foil centre and with the radiation field characterised by the radiation temperature in figure 5.6.5(a).
Figure 5.6.6(a). Same as 5.6.4(a) but at the foil rear.
Figure 5.6.6(b). Same as 5.6.4(b) but at the foil rear and with the radiation field characterised by the radiation temperature in figure 5.6.6(a).
This is contrary to the calculation which predicts that \( \delta(1) \) decreases towards zero. Therefore, the small departure from LTE seen at the rear of the foil is almost entirely a time dependent effect.

In contrast, when a radiation field is included, departures from LTE are seen so that \( \delta(1) < 0 \) during the heating phase. This is because \( T_R > T_e \). Therefore, the radiation will tend to force the material towards a state which is characteristic of a radiation temperature \( T_R \). However, following the discussion in 5.6.1, the populations of the higher lying levels are still forced towards a state reflecting \( T_e \) because collisional processes are more important for these levels. The net result is that \( n = 1 \) is depleted by photopumping of electrons into \( n = 2 \) (and \( n = 3 \) and the continuum) and the populations of \( n = 2 \) and \( n = 3 \) are enhanced only slightly.

At the front of the foil, during the first 200 ps, \( T_R \) rises significantly faster than \( T_e \). This is reflected in the \( \delta(1) \) profile which decreases sharply to just under -1% during this time. From the calculation without a radiation field, it is seen that significant departures from LTE \( \gtrsim 1\% \) begin at around 200 ps. This is responsible for the rapid decrease in the magnitude of \( \delta(1) \) at around 200 ps when the radiation field is included. Up to \( t = 1000 \) ps, \( T_R \approx T_e \) and the departure from LTE is small. Note that it is the presence of the radiation field that maintains the system close to LTE. After this time, \( T_R \) falls off markedly in comparison to \( T_e \) and this is accompanied by a rapid rise in \( \delta(1) \) to around 7%. Note that, at late times, \( \delta(1) \) (with radiation field) < \( \delta(1) \) (without radiation field).

Very similar effects are observed deeper in the target, except that the
departure from LTE is larger when negative than at the front (-6% centre, -30% rear). This is because the differences between $T_R$ and $T_e$ are greater, which can be seen clearly from the temperature plots. Also note that at the back of the target, $s(1) \leq 0$ always. This is because $T_R > T_e$ at all times and the departure from LTE when no radiation field is present is very small.

From these calculations, it is clear that the radiation field has a significant effect on the atomic level populations. Most importantly, these are maintained closer to LTE than when no radiation field is present. The exception is towards the rear of the foil where the radiation temperature becomes much larger than the material temperature. However, the -30% departure in $P_1$ is probably an overestimate. This is because the radiation field close to resonance for $1 \rightarrow 2$ transitions will be less than that described by the values of $T_R$ used in the calculation. This results from the significant "hole" in the spectrum predicted by the LTE radiation transport model during the pulse rise, which is observed at all times in the experiment (5.3). In conclusion, the use of the LTE model in the radiation transport calculation appears to be a very good approximation.
CHAPTER 6

MODELLING OF ULTRASHORT LASER PULSE INTERACTION WITH SOLID TARGETS USING A HYDRODYNAMICS CODE

In this chapter, the modelling of ultrashort laser pulse interaction with solid targets using the 1D Lagrangian hydrodynamics code MEDUSA is described. The case when a significant prepulse produces a preformed plasma with which the high irradiance short pulse interacts strongly is also investigated. This work was performed by the author in support of an experiment in which he was involved (Willi et al 1989). This experiment produced the first observations of a hot (400 eV), high density \( n_e \gtrsim 10^{23} \text{ cm}^{-3} \), fully ionized aluminium plasma formed by the interaction of an ultrashort laser pulse with a solid target. Although the author was mainly involved in performing the hydrodynamics modelling, a description of the experiment and experimental results is presented. The analysis leading to the conclusion that a plasma with an electron density of around \( 1.6 \times 10^{23} \text{ cm}^{-3} \) at a temperature of 400 eV was observed is briefly described also.

Ultrashort pulse interactions with preformed plasmas have been reported recently by other authors (Murnane et al 1989, Cobble et al 1989).

6.1 DESCRIPTION OF EXPERIMENT

The experiment was carried out at the SERC Central Laser Facility (Rutherford Appleton Laboratory) using the KrF short pulse, high
power Sprite laser system (Barr et al 1988). The short pulse was generated by a mode-locked dye laser which operates at a wavelength of 0.746 μm and is amplified in a three stage dye amplifier system. It is then frequency tripled in two KDP crystals resulting in a wavelength of 0.248 μm. Finally, the 3.5 ps pulse is amplified to energies of up to 2.5 J in one discharge and two electron beam pumped KrF amplifiers. The laser beam was focused on target with an f/2.5 aspheric lens to a 20 μm focal spot. The focal spot size was measured with an X-ray pinhole camera with 15 μm resolution and filtered for photons in the kilovolt energy region. A fast optical diode was used to measure the ASE prepulse level. The total energy was recorded with an energy calorimeter.

The main diagnostic used was a time resolved X-ray crystal spectrometer. A flat TLAP crystal was coupled to an X-ray streak camera to operate in the 5 - 7 Å spectral wavelength region. The spectral and temporal resolutions were 15 mA and 15 ps respectively. Time integrated spectra were recorded with a PET crystal. Figure 6.1.1 shows the experimental arrangement.

6.2 EXPERIMENTAL RESULTS

6.2.1 Measurement of Prepulse Level

Energy levels of less than 1% were measured for low prepulse shots using the optical diode. This resulted in a contrast ratio of less than 10⁻⁵ in power since the ASE was emitted during a 20 ns interval. In addition, targets irradiated with large prepulse levels showed craters with diameters in excess of 200 μm in contrast to the 20 μm focal spot produced for the main pulse. This indicates a larger divergence for
Figure 6.1.1. Schematic of experimental arrangement for ultra short pulse solid target interaction experiment.
the ASE which reduces the contrast ratio between prepulse and main pulse still further. However, this is not a quantitative measurement because lateral heat flow may have occurred during the 20 ns ASE pulse duration.

The VUV emission from the targets was also measured using a time resolved VUV spectrometer. For large prepulses containing 10% or more of the total energy, Li - and Be - like transitions of aluminium in the VUV were clearly observable. In contrast, no emission was seen in front of the 3.5 ps laser pulse.

Finally, titanium targets overcoated with layers of plastic, 0.5 μm to 4 μm thick, were used to study the X-ray emission from titanium. A large Kα signal as well as the He - like resonance spectrum was seen for large prepulses, indicating that several microns of plastic were ablated by the prepulse. When the prepulse was small, only the Kα emission was seen on targets with plastic coatings less than 2μm thick, demonstrating that the low level of ASE did not damage the target.

6.2.2 High Density Observations
A number of features associated with high density plasmas such as Stark broadening and the effects of continuum lowering were seen in X-ray spectra from shots when the ASE prepulse level was very low. Figure 6.2.1 shows an X-ray streak spectrum from an Al foil target irradiated with a 3.5 ps, 3 \(10^{16}\) Wcm\(^{-2}\) laser pulse. The spectrum is dominated by the wide, intense He-like 1s\(^2\) → 1s3p transition. Higher series transitions from He - like ions and the Ly8 transition are faint and turn on some 15 ps after the bright He - like line. In contrast, a distinctly different spectrum is seen from shots with a large prepulse.
Figure 6.2.1. X-ray streaked spectrum from an aluminium foil irradiated by a 3.5 ps, $3 \times 10^{16} \text{ Wcm}^{-2}$, KrF laser pulse when the perpulse level was insufficient to cause breakdown of the target surface.

Figure 6.2.2. X-ray streaked spectrum from an aluminium target irradiated by a 3.5 ps KrF laser pulse superimposed on a large ASE prepulse.
Figure 6.2.2 shows a spectrum taken on an Al target irradiated by a pulse with an ASE prepulse containing 30% of the total energy. Clearly observable are all the members of the H-like and He-like series. Also, the width of the He(ls² → ls3p) line is significantly less and bright emission is seen over several hundred picoseconds which indicates the presence of a large plasma corona.

The electron density was obtained by comparing the Stark profiles of the He-like ls² → ls3p and ls² → ls4p transitions with profiles predicted by the atomic physics codes RATION and SPECTRA (Lee, Whitten and Stout 1984). This analysis was carried out by V. Barrow. Figure 6.2.3 shows densitometer traces of the two He-like transitions of figure 6.2.1 taken at approximately 15 ps after the start of the emission. The two line profiles predicted by the atomic physics codes for an electron density of \(1.6 \times 10^{23} \, \text{cm}^{-3}\) are superimposed. The calculation assumed an electron temperature of 400 eV and included opacity corrections assuming a homogeneous plasma slab 1 μm thick. The corrections due to opacity were very small, not affecting the predicted profile significantly for plasma thicknesses up to 1 μm.

A further indication of the high plasma density is the absence of higher series members in the X-ray spectra produced when the ASE prepulse level was low. High density effects in general have been discussed in more detail in 4.7.3.

The temporal evolution of the spectral lines is illustrated in figure 6.2.4. Microdensitometer traces of the streak image shown in figure 6.2.1 are shown at two different times. Initially, the He(ls² → ls3p) transition is very broad and intense, peaking after the start of the
Figure 6.2.3. Comparison between experimentally measured profiles of aluminium $1s^2 \rightarrow 1s4p$ and $1s^2 \rightarrow 1s3p$ transitions and predicted profiles assuming $T_e = 400$ eV and $n_e = 1.6 \times 10^{23}$ cm$^{-3}$. 
Figure 6.2.4. Densitometer traces across the spectrum shown in figure 6.2.1 taken at 20 ps and 35 ps from the start of the emission.
emission. The width changes in time and is asymmetric towards higher energies. At about 25 ps after the start of the emission, the profile decreases in width and becomes weak at about 45 ps. The $1s^2 \rightarrow 1s4p$ He-like transition turns on after about 15 ps. After about 30 ps, the Ly$\alpha$ transition is observed, coincident with the decrease in width of the $1s^2 \rightarrow 1s3p$ line.

The electron temperature was obtained from the ratio $\text{Al}(\text{Ly}\alpha)/\text{Al}(1s^2 \rightarrow 1s3p)$ of time integrated spectra. A temperature of $400 \pm 50$ eV was found by using small opacity corrections.

The electron density was also obtained for the plasma producing the distinctly different X-ray spectrum shown in figure 6.2.2. The analysing procedure was the same as for the spectrum shown in figure 6.2.1. An electron density of $3 \times 10^{22}$ cm$^{-3}$ gave best agreement with the experimental data. This is consistent with emission from densities close to critical ($1.8 \times 10^{22}$ cm$^{-3}$) for 0.248 $\mu$m laser light.

6.2.3 Ti K$\alpha$ Measurements
Titanium foil targets with various thicknesses of plastic coatings were used to investigate the level of hot electron production in shots with and without a large prepulse. The K$\alpha$ signal was measured with time integrated X-ray crystal spectrometers. Figure 6.2.5 shows microdensitometer traces of K$\alpha$ spectra taken on a bare titanium target and on one coated with 0.5 $\mu$m of plastic. No signal was observed when the plastic coating was 2 $\mu$m thick and the prepulse level was very low. From the K$\alpha$ signal of these shots, it was found that 20% of the incident laser energy was deposited in hot electrons with energies larger than 4.5 keV. In contrast, when the prepulse level was
Figure 6.2.5. Densitometer traces of $K_{\alpha}$ spectra taken on a bare and 0.5 $\mu$m plastic coated titanium foil target.
large, a strong Kα signal was observed together with He-like titanium lines, indicating that several microns of plastic were ablated by the prepulse.

6.3 HYDRODYNAMICS CODE SIMULATIONS

Extensive numerical modelling was performed using the 1D Lagrangian hydrodynamics model, MEDUSA, without knowledge of the absorbed laser energy. All the simulations were carried out with a laser wavelength of 0.248 μm and aluminium and plastic coated titanium as the target material. The low prepulse calculations were made with a 3.5 ps (FWHM) Gaussian pulse containing 0.55 J to model the data shown in figure 6.2.2. To simulate the conditions which produced the spectrum shown in figure 6.2.3, a 3.5 ps pulse containing 0.55 J was superimposed on a 20 ns background pulse containing 0.25 J (~ 30% of the total energy).

In the modelling, the laser energy is absorbed via inverse bremsstrahlung up to the critical surface, where a fraction $f_a$ of the remaining energy is dumped. A fraction $f_h$ of the dumped energy is assumed to produce hot electrons via resonance absorption. The hot electron temperature is obtained from an expression determined experimentally (Giovanelli et al 1976). The resulting energy spectrum of the suprathermal electrons is split into 10 groups. These are transported through the target with a simplified model (Evans 1980) which uses the energy loss equation for fast test particles as given in NRL Plasma Formulary. The energy remaining in hot electrons when one pass through the target has been made is distributed equally
throughout the target. For most purposes, as is the case here, the targets are sufficiently thick that the energy remaining in the fast electrons after one target pass is negligible.

6.3.1 Simulations Without Prepulse
Several simulations were carried out varying $f_a$ between 0.01 and 0.5 (the theoretical maximum (Hughes 1979)) because the absorbed fraction of laser energy was not known at the time. Also, because the fraction of anomalously absorbed energy producing hot electrons was not known, $f_h$ was varied between 0.1 and 0.9. A larger value of $f_h$ was expected to be more realistic. The classical electron flux limiter $f_c$ was set to 0.1 on the basis that the laser plasma interaction region, $d$ was much less than the focal spot size, $D$, so that lateral heat flow could be expected to be small in this region. $d$ was estimated from $d \approx c_s \tau_{\text{FWHM}}$, where $c_s$ is the plasma sound speed and $\tau_{\text{FWHM}}$ the laser pulse width. Even for a plasma temperature of 1 keV, $d \approx 1 \mu\text{m} \ll D(\approx 20 \mu\text{m})$. This choice of $f_c$ was later verified by Rickard, Bell and Epperlein (1989). Simulations were also carried out for $f_c = 0.03$, the common value for 1D hydrocodes modelling planar targets. Large differences in the results were only seen, however, when $f_h$ was very small.

The best agreement with the experimental data was found for $f_a = 0.2$ and $f_h = 0.9$ i.e. 18% of absorbed energy producing hot electrons. In figure 6.3.1, MEDUSA output, 2 ps after the peak and at the end of the pulse, is shown for the electron density, electron temperature and average degree of ionization. At the earlier time, an electron temperature of just over 600 eV is predicted for $n_e \sim 10^{23} \text{ cm}^{-3}$.
Figure 6.3.1. Hydrocode predictions of plasma conditions 2 ps (top) and 4 ps (bottom) after the peak of a 3.5 ps laser pulse assuming 20% of the laser energy reaching critical density to be anomalously absorbed.
Figure 6.3.2. Same as figure 6.3.1 but assuming an anomalous absorption fraction of 10%.
Figure 6.3.3. Same as for figure 6.3.1 but assuming an anomalous absorption fraction of 30%.
By the end of the pulse (4 ps after the peak), temperatures between 400 eV and 500 eV are observed for $n_e \sim 2 \times 10^{23}$ cm$^{-3}$, in close agreement with the experimental observations. For comparison, in figure 6.3.2 and 6.3.3, the predicted profiles are shown at the same simulation times for $f_a = 0.1$ and 0.3 respectively. All other parameters are the same as for the calculation giving the results shown in figure 6.3.1. As can be seen, when $f_a = 0.1$, the plasma remains comparatively cool, not exceeding 370 eV at electron densities over $10^{23}$ cm$^{-3}$. In contrast, when $f_a = 0.3$, electron temperatures close to 1 keV are predicted at an electron density of $10^{23}$ cm$^{-3}$. Clearly, the predicted plasma conditions are strongly dependent upon $f_a$, as may be expected.

Although results are presented above for just three different values of a single parameter, values of $f_a = 0.2$ and $f_h = 0.9$ were the only combination of parameters of the very many assessed to produce a plasma closely resembling that which was experimentally observed. Note that although densities close to solid are predicted during the laser pulse, this is not observed experimentally because of the finite temporal resolution of the streak camera (15 ps).

In common with most hydrodynamics modelling (1D and 2D), the plasma was found to cool too rapidly. In figures 6.3.4, 6.3.5 and 6.3.6, plasma profiles are presented at 15 ps and 50 ps after the peak of the pulse for $f_e = 0.9$ and $f_a = 0.1$, $f_a = 0.2$ and $f_a = 0.3$ respectively. At 15 ps after the peak of the pulse, the plasma density is still quite high ($n_e \approx 3 \times 10^{22}$ cm$^{-3}$) although the temperature has fallen to less than 300 eV (for $f_a = 0.3$). By 50 ps, although the plasma density is predicted to be above $10^{22}$ cm$^{-3}$, the temperature only just exceeds
Figure 6.3.4 Hydrocode predictions of plasma conditions 15 ps (top) and 50 ps (bottom) after the peak of a 3.5 ps laser pulse assuming 10% of the laser energy reaching critical density to be anomalously absorbed.
Figure 6.3.5. Same as figure 6.3.4 but assuming an anomalous absorption fraction of 20%.
Figure 6.3.6. Same as for figure 6.3.4 but assuming an anomalous absorption fraction of 30%.
100 eV for \( f_a = 0.2 \) and \( f_a = 0.3 \) while the temperature has fallen significantly below 100 eV for \( f_a = 0.1 \). In contrast to these predictions, the plasma temperature, 50 ps after the start of the emission, is experimentally measured to be close to 350 eV. On the other hand, the density is inferred to be several times \( 10^{23} \text{ cm}^{-3} \) from comparison of the experimentally measured line profile of the aluminium He - like \( 1s^2 \rightarrow 1s3p \) transition with profiles predicted by RATION and SPECTRA. This is from a preliminary analysis supplied by V. Barrow.

6.3.2. Simulations With Prepulse.

Simulations were also carried out with various levels of prepulse superimposed on the 3.5 ps short pulse to check the consistency of the predictions with the experimental data. In particular, a 20 ns Gaussian pulse containing 30\% of the short pulse energy was superimposed on a 3.5 ps 0.8 J pulse.

The simulations clearly show that during the prepulse several microns of material are ablated, resulting in a large preformed plasma corona. Figure 6.3.7 shows the predicted plasma conditions due to the prepulse just before the arrival of the high intensity short pulse. An electron temperature of around 100 eV is observed in the underdense plasma which is consistent with the VUV emission observed long before the arrival of the main pulse.

The high intensity short pulse was completely absorbed in the large underdense plasma, resulting in temperatures of several keV during the interaction. In contrast to the prepulse free case, a significant classical electron flux limiter of 0.03 had to be used to simulate the
Figure 6.3.7. Hydrocode predictions of plasma conditions due to a low irradiance prepulse incident on an aluminium target.
effect of lateral heat flow, thus preventing the high density plasma from becoming unrealistically hot. This choice of flux limiter was later verified by 2D Fokker-Planck simulations (Rickard, Bell and Epperlein 1989) which showed that the heat flow was largely lateral in the underdense plasma, limiting temperatures to approximately 1 keV.

Figure 6.3.8 shows the predicted plasma conditions just after the high intensity short pulse interaction with the plasma shown in figure 6.3.7. Note the high temperature in the underdense plasma and the rapid fall off on going to higher densities resulting from the 0.03 classical electron flux limiter. Also, several microns of material have been ablated. This is consistent with the observation of He-like resonance transitions of titanium when titanium foil targets, coated with 4 μm of plastic, were irradiated by the main pulse superimposed on a large ASE prepulse.

6.4 PRELIMINARY MODELLING FOR THE IMPERIAL COLLEGE SHORT PULSE LASER SYSTEM.

To attempt to predict the plasma conditions that may be expected to be produced by the 300 fs (FWHM), 1 μm laser system shortly to become operational at Imperial College, preliminary simulations have already been carried out in a similar manner to those described in 6.3.1 when no prepulse is present. As an example, it has been assumed that the anomalously absorbed fraction of the laser energy reaching the critical surface is 20%, 90% of which is assumed to produce hot electrons as before. The hot electron temperature was calculated in the same way described above.
Figure 6.3.8. Hydrocode predictions of plasma conditions just after the interaction of a 3.5 ps, $3 \times 10^{16} \text{ Wcm}^{-2}$, KrF laser pulse with the aluminium plasma whose conditions are shown in figure 6.3.7.
Figure 6.4.1 shows the predicted plasma conditions for an aluminium target close to the end of 300 ps, $10^{17}$ Wcm$^{-2}$ laser pulse, which are expected to be routinely delivered to the target surface. A plasma with a temperature above 0.5 keV is predicted to be formed at "solid" electron density ($7.8 \times 10^{23}$ cm$^{-3}$). These types of plasmas will be extremely interesting for dense plasma physics studies and research in stellar atmospheres.

It should be noted that these preliminary results should be regarded as such. They require further investigation, particularly because the laser absorption mechanisms are unclear and the atomic physics in the model is probably not adequately treated at such short pulse lengths. It is these areas that must receive most attention in future work.
Figure 6.3.9. Preliminary hydrocode predictions showing that hot plasma might be produced at solid density by a 300 fs, 1 µm, 10^{17} Wcm^{-2} laser pulse. The laser absorption fraction is 20%, most of which is assumed to produce hot electrons.
CHAPTER 7

CONCLUSIONS

The author has designed and performed an experiment to measure radiation transport through gold coated thin foil plastic targets of various thicknesses using high resolution time resolved XUV spectroscopy in the 10 A to 70 A spectral wavelength region for the first time. The results have been predicted and analysed with a multigroup radiation transport model using 116 energy groups in the 0 to 100 keV energy region and an in-line opacity calculation, assuming LTE, which was developed by the author. Good overall agreement was obtained between the predicted and measured spectra emitted from the rear of the targets. In particular, the edge feature, the progressively retarded emission towards longer wavelengths and the early turn on of radiation just below (in energy) and well above the edge are reproduced by the simulations.

The author's extensive modelling of the radiatively heated foils has allowed him to draw the following important conclusions. The foil heating is a radiation transport process, even though the material is made from low Z constituents. The energy transfer in the foils is completely dominated by radiation transport. Edge shifting, of the carbon K-edge in particular, contributes considerably to the propagation of the heat wave. Edge smearing, not previously considered, results from the temperature gradient generated in the sample and gives rise to a smeared out opacity spectrum, altering the power deposition profiles in the material. Care should be employed when using tabulated opacity data to ensure that edge shifting and smearing are
It has been clearly demonstrated that it is the radiation field that maintains much of the plasma close to LTE. Previously, the validity of the assumption of LTE has been considered on the basis that collisional processes must be dominant. It was also shown that at the rear of the targets, the radiation field may drive the plasma significantly away from LTE.

Some discrepancies between the experimental results and predictions were found. The small blue shift of 1 Å in the edge feature on the experimentally measured spectrum from a 5 μm target is not reproduced by the model, which predicts a blue shift in the carbon K-edge of more than 10 Å. It has been shown that this is probably due to the filling in of the opacity spectrum below the shifted K-edge by strong carbon absorption lines not considered in the model.

It has been demonstrated that differences between the predicted and experimentally measured temporal profiles of the XUV emission, away from the edge, from the rear of the targets is probably due to differences in the assumed temporal profile of the heating radiation from the gold and/or variations in the laser pulse length or shape. However, this cannot explain the significant delay in the emission of radiation just above (in energy) the cold K-edge position, predicted by the code and not experimentally observed. It was shown, however, that local variations in the spectral or temporal shape of the radiation pulse do not affect the overall results or conclusions.

The author has performed extensive hydrocode simulations of ultrashort
(3.5 ps), high intensity ($\approx 10^{16}$ W cm$^{-2}$), KrF laser pulse interaction with solid targets, in support of an experiment in which he was involved and which produced the first observations of hot (400 eV) high electron density ($\approx 10^{23}$ cm$^{-3}$) plasma formed in this way. Best agreement with the experimental observations was obtained when it was assumed that approximately 20% of the incident laser energy was absorbed and converted into hot electron energy and a classical electron flux limiter of 0.1 was used. The assumed absorption percentage was later verified by analysis of measurements of K$_\alpha$ signals from bare and plastic coated titanium targets.

Simulations with a large prepulse required a significant flux limiter of 0.03. This was to prevent the high density plasma from becoming unrealistically hot due to the plasma corona produced by the prepulse, reaching temperatures of several keV when it completely absorbed the main pulse. In reality lateral heat flow prevents such high coronal temperatures.

Finally, the experience gained from the short pulse modelling has been used to predict that above 0.5 keV plasmas at solid density may be produced by the Imperial College 300 fs high power laser system currently under construction. These are preliminary results because absorption is uncertain and ionization physics are probably not adequately treated by the code on ultrashort time scales.
REFERENCES


