Laser-plasma coupling effects on spectral line shapes

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Abstract

This thesis presents an experimental and computational study of strong electric field effects on plasma spectral line shapes. The role of laser-driven electron plasma waves is studied for helium-like systems where atomic energy levels are weakly susceptible to external electric fields. In this case strong electric fields associated with laser driven electron plasma waves are resonant with the radiators producing modulations on line emission. In an experiment the ASTRA laser system was used to generate expanding plasma ionised to H-like and He-like states using short (picosecond) laser pulses. Spectral emission dispersed by a high dispersion toroidal spectrometer capable of high spatial and spectral resolutions is presented. Cross sections of the emission close to the target surface contains modulations on the He $\beta$ and He $\gamma$ resonance lines. The modulations are interpreted using the Intra-Stark framework. The modulations in each data set were found to be emitted from and below quarter critical density. By introducing a femtosecond pulse after the picosecond pulse the modulations were reduced and, in some cases, suppressed altogether. Plasma hydrodynamics and laser-plasma coupling processes are simulated using a hydrodynamics code (EHYBRID) and post-processed to include parametric processes. Simulations suggest that high amplitude electron plasma waves are generated in the plasma underdense to the laser. Atomic kinetics and spectral lines are modelled using the CRETIN code to examine the temporal evolution of the resonance line emission. Simulated spectra are post-processed to include the effects of strong high frequency electric fields associated with the laser driven electron plasma waves. The simulated spectra reproduce the modulations.
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Author’s declaration

The work presented in this thesis is, unless stated otherwise, the work of the author and has not previously been submitted for a degree at this or any other university. The experimental work is part of a collaborative effort with a large number of people at large-scale experimental facilities and other institutions. I have been involved in all aspects of the research process. This includes experimental design, diagnostic characterisation, execution of experiments, analysis of the results and publication.

The setup and maintenance of the large area cooled-chip CCD and assistance in setting up the high dispersion toroidal crystal spectrometer was undertaken by the author. Computer simulations using the EHYBRID and CRETIN codes in chapters 5 and 6 are the work of the author. EHYBRID and CRETIN post-processors were solely developed by the author. Portions of the work presented here have been published or have been accepted for publication [1–5].
Chapter 1

Introduction

This chapter presents an introduction to the subject matter and motivation of the thesis. Electric fields in plasma are introduced and their ability to accelerate particles to high energy. Plasma spectroscopy is then presented, motivated by the need for diagnosing densities and temperatures within a plasma. The effect of strong electric fields on spectral transitions and their spectroscopic signature is discussed. Finally experimental considerations are presented.

1.1 Electric fields in plasma

The measurement of electric fields in laser produced plasma is of particular importance. Large electric fields can accelerate particles into collimated high energy beams for uses such as material probing [6–8]. This can be achieved on a much smaller scale than conventional radio-frequency cavities such as CERN, limited by electric breakdown to $\sim 10^6$ V cm$^{-1}$. In contrast these high energy particles can be problematic. Inertial confinement fusion (ICF) is a possible route in generating energy using controlled thermonuclear explosions [9–12]. Here a spherical target containing deuterium-tritium (DT) is irradiated uniformly across it’s sur-
face. The requirement is to compress the DT fuel to very high densities (with $\rho r$ a product of density and radius of the compressed fuel). Once the fuel is compressed it is ignited either by a central hot spot or by a separate fast ignition laser. The temperatures required are $\sim 10$ keV. The combination of high densities and high temperatures are required to raise the tunnelling probability and achieve energy gain. The $\rho r$ product is required to trap $\alpha$-particles and drive a fusion burn throughout the fuel. What is critical is that the DT fuel remains cool during the compression phase. It is essential that the laser does not accelerate particles which may preheat the fuel ahead of the compression.

The electric fields that generate high energy particles are associated with laser driven electron plasma waves and are amplified by the interactions between the electromagnetic waves associated with the laser beam and the plasma. The electric fields associated with the electron plasma waves can reach values greater than $10^9$ V cm$^{-1}$ when amplified by laser interaction, dependent on the electron density and temperature. Laser-plasma coupling processes are well known and have been studied extensively [13–19]. However mechanisms where the energy of the laser is completely absorbed into driving electron plasma waves, for example two-plasmon decay, can only be observed with probes such as Thomson scattering or particle emission [20]. In this case only the presence of the high amplitude electron plasma waves can be detected.

1.2 Plasma spectroscopy

The measurement of plasma properties can be made by studying the emission from bound-bound, bound-free and free-free transitions of the electrons (discussed in chapter 3). For hot, high density laser-produced plasma it is sometimes the only method available. Conventional probes (such as Langmuir probes) can affect the properties of the plasma. The emission from bound-bound transitions is
directly influenced by the plasma temperature and density. Plasma spectra therefore contain information about hydrodynamic parameters and atomic processes. Radiating ions are also affected by strong electric fields. Low frequency fields (with respect to the duration of the transition) associated with the ion charge shift the energy levels of the transition. The line emission from bound-bound spectra is broadened by this process. The transition can be perturbed by high frequency electric fields associated with electron plasma waves and manifest as modulations (or plasma satellites) in emitted spectra. The position and structure of the modulations is indicative of the frequency and electric field strength of electron plasma waves. Similarly the transitions can be perturbed by an external high frequency field such as that associated with a high intensity laser (laser satellites).

The theory of plasma satellites was introduced by Baranger and Mozer in 1961 and developed by a number of groups [21–25]. The action of the field on the upper atomic level creates additional channels of transition in helium-like emission. Here the modulations consisted of local maxima where intensity abruptly increases. Initial measurements used hollow-cathode arc discharge or z-pinches to create low density low temperature plasma [26]. In this case high frequency electric fields were generated by capacitor discharge into the plasma. In high intensity laser plasma experiments plasma emission is mainly composed of high energy photons (X-rays). For higher values of atomic number $Z$ bound-bound photon energies scale as $Z^2$. Observations of modulations are particularly difficult due to the high resolving power needed at these photon energies. The development of bent geometry X-ray diffracting crystals with high spectral dispersion has allowed research groups to study these modulations in laser-plasma experiments [27]. Time resolved measurements of satellites on the Al$^{+11}$ He-like emission has been reported [28]. In this case the electron density calculated from the satellite positions is supported by hydrodynamic and atomic kinetic simulations.
Zhuzhunashvili and Oks initially studied the role of a resonance between the Stark split levels of bound-bound transitions and electron plasma waves [29]. Here the frequency of the electron plasma wave is equal to the Stark shift of the atomic levels. This effect is particularly strong in hydrogen-like transitions where the large Stark splitting occurs for low electric field strengths. Modulations in the spectra consist of local maxima and minima. The electric field strength can be directly measured from the structure of the modulations. This theory was developed for high density laser produced plasma and supported a number of papers [30–33]. Z-pinch and laser-plasma experiments have verified the existence of these modulations in hydrogen-like lines [34–36]. Hydrodynamic and atomic kinetic simulations have supported the measurements. It has been reported that spectroscopic measurements of ICF targets have contained these modulations [30,37–39].

A relatively unexplored area is the action of resonance on helium-like transitions. Here the low frequency ion fields must be large enough to split the atomic transition levels substantially. The high frequency fields associated with the electron plasma waves are resonant with this ensemble of atoms. In addition the high frequency electric fields must be large enough for resonance to occur. Modulations on the line shape will occur for large amplitude laser driven electron plasma waves. By studying the modulations produced by this resonance it is possible to infer the presence of laser-electron plasma wave coupling, the densities at which it occurs and the electric field strength associated with the waves using spectroscopic methods. This is an innovative approach to study laser-plasma coupling.

In the experiment described in this thesis a highly dispersive imaging crystal spectrometer was used to resolve plasma emission. The plasma was spatially resolved along the target normal. Close to the target surface plasma expansion is approximately 1-dimensional. With this approximation the plasma parameters can be specified using this single coordinate axis. Using a toroidally bent
spectrometer it was possible to resolve approximately 0.1 eV in 1800 eV. This condition is necessary to observe modulations due to laser driven electron plasma waves. By analysing the plasma emission from different spatial locations the hydrodynamic and atomic kinetic parameters can be inferred as a function of height above the target surface. 1-dimensional hydrodynamic and atomic kinetic simulations of the plasma parameters can be compared to data. Emission was not temporally resolved. Time resolved measurements are valuable but technical difficulties limit spectral resolution. Furthermore high spatial resolution cannot be achieved when temporally resolving emission.

1.3 Chapter outline

Chapter 2 (Laser-plasma coupling) In this chapter the theory of laser-plasma coupling is introduced. The dispersion relations for electromagnetic and electron plasma waves are presented. Laser absorption by collisional and non-collisional (wave) mechanisms is discussed. The amplification of electron plasma waves by laser-absorption is investigated. Wave damping and suppression is evaluated. Laser generated magnetic fields are presented.

Chapter 3 (Plasma spectroscopy) Plasma spectroscopy is discussed with special reference to Stark splitting. The quadratic and linear Stark effect in helium-like systems is calculated. Dynamic resonance between the electron plasma wave and Stark split energy levels for high density laser produced plasma is discussed.

Chapter 4 (The ASTRA experiment) The ASTRA experimental investigation of laser-plasma coupling is presented. This includes the experimental setup, target design and diagnostics. Cross sections of the spectral emission are examined and compared to theory.

Chapter 5 (Hydrodynamics simulations) In this chapter the hydrodynamic
simulations of the experimental plasma using the EHYBRID code are presented. Output from the EHYBRID code is post-processed to calculate the presence, growth and damping of electron plasma waves. These are compared with experimental measurements.

Chapter 6 (Lineshape simulations) Simulated spectra are produced by the atomic kinetics code CRETIN using the EHYBRID output. These simulated spectra are compared to data to support the accuracy of the simulations. Spectra are post-processed to include modulations created by dynamic resonance. The position and structure of the modulations are compared to data.

Chapter 7 (Conclusions) Conclusions and suggested further work is presented.
Chapter 2

Laser-Plasma Coupling

2.1 Introduction

In this chapter the theory of laser-plasma coupling is introduced. Laser absorption by collisional and non-collisional mechanisms is described. Plasma waves supported by plasma are discussed using fluid and kinetic treatments. These can be grouped into three categories; electromagnetic waves, and electron, and ion plasma waves. Due to the short time scales under investigation and their large mass the motion of the ions can be neglected. Amplified waves driven by laser absorption are presented. This includes resonance absorption and parametric instabilities which occur at different locations within the plasma. Self-focusing mechanisms are discussed where the laser alters the refractive index of the plasma and creates filamentation. Finally magnetic fields generated by the laser are presented.

Figure 2.1 illustrates the spatial profile of electron densities in an inhomogeneous plasma. Figure 2.1 shows the plasma density approximated as an exponential decrease, the locations of laser plasma coupling mechanisms are highlighted. Resonance absorption occurs at or close to the critical surface where electron
densities $n_e$ are equal to the critical electron density $n_{cr}$. The critical electron density is defined as the density $n_e$ at which the laser frequency $\omega_L$ is equal to the electron plasma frequency $\omega_{pe}$. Laser driven parametric instability mechanisms take place in the lower density plasma where $n_e < n_{cr}$. Laser photons propagating through the low density plasma are scattered into photons and plasmons, the quantum equivalent of a plasma wave. The process can feedback and cause amplification of plasmons. The location of parametric instabilities is dependent on the process. Stimulated Brillouin scattering (SBS) takes place at or below $n_{cr}$, and Stimulated Raman Scattering (SRS) and Two-Plasmon Decay (TPD) occur at densities where $n_e \leq n_{cr}/4$ or quarter critical density.

In this chapter the dispersion of electromagnetic waves is discussed regarding the nature of laser light travelling through a plasma. Electron plasma oscillations are discussed using fluid and kinetic treatments for both “cold” ($T = 0$) and “hot” ($T > 0$) conditions. The importance of noncollisional damping of the waves and the creation of suprathermal electrons is described. Finally, the interaction between different wave modes resulting in wave amplification is discussed. This includes resonance absorption and the parametric instabilities SRS and TPD.
2.2 Electromagnetic waves

To analyse the motion of particles within a plasma due to the presence of transverse and/or longitudinal waves the plasma is treated as a volume of fluid containing individual fluid elements. Each element contains a mass \( m_s n_s dV \) and charge \( q_s n_s dV \) where \( m_s, q_s, n_s \) is the mass, charge, number density of particle species \( s \) (electrons or protons) and the volume element they occupy respectively. Each element will be subject to the Lorentz force \( \mathbf{F}_L = q_s n_s dV \left( \mathbf{E} + \mathbf{v}_s \times \mathbf{B} \right) \) and a force due to pressure gradients between each element \( \mathbf{F}_P = \nabla P_s dV \). \( \mathbf{E} \) and \( \mathbf{B} \) are the electric and magnetic fields experienced by the particles, \( \mathbf{v}_s \) is the velocity of species \( s \) and \( P_s \) is the pressure of each species. The equation of motion can now be written using the convective derivative \( d/dt = \partial/\partial t + \mathbf{v}_s \cdot \nabla \) as:

\[
\frac{\partial \mathbf{v}_s}{\partial t} + \mathbf{v}_s \cdot \nabla \mathbf{v}_s = \frac{q_s}{m_s} \left( \mathbf{E} + \mathbf{v}_s \times \mathbf{B} \right) - \frac{1}{m_s n_s} \nabla P_s \tag{2.1}
\]

The flow of charge due to the Lorentz and pressure gradient forces on the fluid elements is described using the current density \( \mathbf{J} \). The ions are considered to be motionless due to their large mass and not effected by an electromagnetic field. Only the electrons contribute to the current density, therefore \( \mathbf{J} = -e n_e \mathbf{v}_e \) where \( e \) is the electron charge. From Maxwell’s equations,

\[
\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad \nabla \times \mathbf{B} = \frac{1}{c^2} \frac{\partial \mathbf{E}}{\partial t} + \mu_0 \mathbf{J} \tag{2.2}
\]

By eliminating \( \mathbf{B} \) from the Fourier analysis of 2.2 we obtain the dispersion relation for electromagnetic waves in unmagnetized plasma,

\[
\omega^2 = \omega_{pe}^2 + k^2 c^2 \tag{2.3}
\]

Here \( \omega_{pe} = \sqrt{n_0 e^2 / m_e \epsilon_0} \) is the electron plasma frequency and \( k \) is the wave...
number of the electromagnetic wave. Figure 2.2 presents the dispersion relation for electromagnetic waves. If electrons are displaced from their charge-neutral state they will collectively oscillate about the original position due to the restoring force caused by $E$ between the separated charged particles with a frequency equal to the electron plasma frequency. This is the only electron frequency permitted in a cold plasma. Therefore the group velocity $v_g = d\omega/dk = 0$ suggesting that a wave packet made of plasma oscillations would not propagate.

As an electromagnetic wave travels through an inhomogeneous plasma of increasing density (increasing $\omega_{pe}$) it’s wavenumber will decrease until it becomes zero at $\omega = \omega_{pe}$ and cannot propagate any further (any solutions where $\omega < \omega_{pe}$ would contain an imaginary component). The wave travels through the under-dense plasma, a region where $\omega > \omega_{pe}$, until it reaches the critical surface where $\omega = \omega_{pe}$. This cutoff occurs at a critical density $n_{cr} = \epsilon_0 m_e \omega^2/e^2$.

### 2.3 Electron plasma waves

In this section the derivation of electron plasma waves is presented. For warm plasmas where $T > 0$ thermal pressure gradients allow plasma waves to propagate. Firstly a fluid treatment of the plasma is discussed and used to calculate the
plasma wave dispersion. A kinetic treatment is then used to analyse the effects of wave damping.

2.3.1 Fluid treatment

For warm plasmas the thermal motion of the particles generates pressure gradients and produces thermal flow between neighbouring volume elements. This thermal pressure gives rise to density compression and rarefaction thereby allowing the plasma oscillations to travel as a wave. In the quantum mechanical picture if we consider an electron plasma wave as a wave packet the frequencies of oscillation have an associated energy $\hbar \omega$. This wave packet, or plasmon, is analogous to the electromagnetic photon. It is a collective oscillation of the electrons. The extra pressure term $\nabla P_s$ in the equation of motion [2.1] must be included to describe the dispersion relation of electron plasma waves (Langmuir wave). This can be found by generating an equation of state which describes the pressure gradients in terms of the thermal velocity $v_{te}$, mass and number density of the electrons within the volume elements. If the thermal conductivity is low the equation of state of the plasma may be written in the form $P_s = C n_s^\gamma$ where $C$ is a constant and $\gamma = c_P/c_V$ is the ratio of specific heats at constant pressure and volume. For a process that contains $N$ degrees of freedom then $\gamma = (N+2)/N$. In extremely fast compressions, such as laser-target irradiation, the compression is adiabatic and, as electron plasma waves are longitudinal in nature, only one degree of translational freedom is found. Therefore the ratio of the specific heats $\gamma = 3$. Using the equation of state in conjunction with the ideal gas equation $P_e = n_e k_B T_e$, where $k_B$ is Boltzmann’s constant and $T_e$ is the electron temperature, to find the last term in the equation of motion [2.1];

$$\frac{\nabla P_s}{m_s n_s} = \gamma \frac{k_B T_s}{m_s} \frac{\nabla n_s}{n_s} = v_{te}^2 \frac{\nabla n_s}{n_s}$$

(2.4)
2.3 Laser-Plasma Coupling

Figure 2.3: Dispersion relation for electron plasma waves. $\omega$ (solid line) tends to the gradient $\sqrt{3}v_{te}$ (dashed line) as $k$ increases.

Here $v_{te} = \sqrt{3k_B T_e/m_e}$ is the electron thermal velocity. By linearising the first order equations and including this extra term we find the dispersion relation for electron plasma waves, also known as the Bohm-Gross relationship [40];

$$\omega_{ck}^2 = \omega_{pe}^2 + 3k^2 v_{te}^2$$  \hspace{1cm} (2.5)

Figure 2.3 shows the Bohm-Gross relationship. As the temperature of the plasma rises from $T = 0$ an increasing spectrum of frequency modes can be supported by the plasma. This is particularly significant as this allows instabilities with differing frequencies to form and grow in hot plasma conditions.

The derivation above assumes that the plasma contains a Maxwellian distribution of particle velocities, obtained by averaging over the distribution functions. It does offer some very important information about electron plasma waves and the frequencies at which they oscillate. However when the same problem is approached using kinetic equations the non-Maxwellian portion of the distribution of the particles introduces extra terms. These extra terms describe the damping of plasma waves.
2.3 Laser-Plasma Coupling

2.3.2 Kinetic treatment

To develop a kinetic theory firstly we consider the plasma to be a collection of non-interacting particles which can be described by the distribution function \( f(x, v, t) \). If, in a collisionless plasma, \( f(x, v, t)d^3x d^3v \) is the number of particles, at time \( t \), in the volume of phase space \( d^3x d^3v \) between \( x \) to \( x + dx \) and \( v \) to \( v + dv \) then the Boltzmann equation must be satisfied;

\[
\frac{\partial f}{\partial t} + v \cdot \frac{\partial f}{\partial x} + F \cdot \frac{\partial f}{\partial v} = 0 \tag{2.6}
\]

which states that the flux of particles into and out of a volume element must be equal i.e. there must be a conservation of particles in the system. By replacing the force term \( F \) with the Lorentz force and ignoring \( B \) the Vlasov equation is formed. For a one-dimensional system;

\[
\frac{\partial f}{\partial t} + v s \frac{\partial f}{\partial x} + q s E \frac{\partial f}{m s \partial v} = 0 \tag{2.7}
\]

When a wave propagates through the plasma the local distribution function will be perturbed by an amount \( \tilde{f}(x, v, t) \) due to the field \( \tilde{E} \) where the \( \sim \) sign signifies a small perturbation. The new distribution function can be written \( f = f_0 + \tilde{f} \). By substituting these values into the Vlasov equation, linearising to first order and Fourier transforming (assuming the perturbations are made up of plane waves of the form \( \exp(ikx - i\omega t) \)) then for electrons;

\[
\tilde{f} = -\frac{q \tilde{E}}{m} \frac{\partial f_0}{\partial v} \frac{i(kv - \omega)}{i(kv - \omega)} \tag{2.8}
\]

By integrating \( \tilde{f} \) over all velocities we can find the number density perturbation \( \tilde{n} \). Using Poisson’s equation \( \partial \tilde{E}/\partial x = \tilde{\rho}/\epsilon_0 \) and letting \( \tilde{E} = -\partial \tilde{\phi}/\partial x = -ik\tilde{\phi}, \) from which \( k^2\tilde{\phi} = \tilde{n}q/\epsilon_0, \) gives the dispersion relation of electron plasma waves.
2.3 Laser-Plasma Coupling

in integral form;

\[ k^2 = \frac{q^2}{m\epsilon_0} \int_{-\infty}^{+\infty} \frac{\partial f_0/\partial v}{v - \omega/k} dv \]  

(2.9)

This integrand contains a singularity at the point \( v = \omega/k \) which can be handled using contour integration in a complex plane. The original prescription was first performed by Landau \[41\], it can be found in standard texts such as \[42\]. The solution of this integral is;

\[ \omega = \omega_{pe} \left[ 1 + \frac{3}{2} k^2 \lambda^2 - \frac{i}{(k\lambda_D)^3} \sqrt{\frac{\pi}{8}} e^{-\frac{1}{2}(k\lambda_D)} \right] \]  

(2.10)

where \( \lambda_D = \sqrt{kT_e/4\pi n_e e^2} \) is the Debye shielding length, a characteristic length beyond which the potential of a charged particle is exponentially attenuated due to the presence of neighbouring charged particles. It is clear that the first two terms in the brackets is the Bohm-Gross dispersion relationship as found in the fluid treatment.

The final term limits the plasma wave to \( k\lambda_D < 0.3 \) \[20\]. The damping of the waves due to non-collisional or dissipative mechanisms is named Landau damping. Physically, this occurs due to electrons being caught in a potential minimum of the plasma perturbing wave. Assuming a Maxwellian distribution, more electrons will be accelerated than are slowed to the phase velocity of the wave \( v_p \). Transfer of energy from the wave to the particles creates a high energy tail in the distribution as shown in 2.4. If the wavelength \( 2\pi/k \) becomes small enough to be comparable to the Debye length \( \lambda_D \) then the effect of Landau damping becomes significant. For \( k\lambda_D = 0.3 \) and above the Landau damping term becomes \( 1/10^{th} \) of the Bohm-Gross term and heavy damping occurs.

Another mechanism that limits the growth of electron plasma waves occurs when the amplitude of the wave becomes so great that the particles along the wave vector \( k \) are disturbed. In this case the wave destroys itself via a process known as wave breaking. For a plasma wave travelling along \( x \) the electrons oscillate with
a displacement $X$ about a mean position $x_0$. Electrons will pass static ions with a charge per unit area of $n_e e X$, where $n_e$ is the undisturbed electron density. An excess positive charge of $n_e e X$ will occur in the $-x$ direction and negative charge of the same magnitude will occur in the $+x$ direction. Using Gauss’s theorem the electric field $E$ on the electrons is:

$$E = \frac{n_e e}{\varepsilon_0} X \quad (2.11)$$

Upon differentiating with respect to time the normal oscillatory motion at the plasma frequency $\omega_{pe}$ is recovered as introduced in section 2.2. If the displacement can be described as

$$X = A \sin(kx_0) \cos(\omega t) \quad (2.12)$$

Where $A$ is the amplitude of the wave. Then at time $t = 0$ in terms of $E$ and $A$
the solution is;

\[ X = \frac{\epsilon_0 E}{n_e e} \]  
\[ \sin(kx_0) = \frac{\epsilon_0 E}{n_e e A} \]  
\[ x = x_0 + X \]

Writing \( E_0 = n_e e A/\epsilon_0 \), and \( E/E_0 = y \) so that \( y \) lies between \( \pm 1 \);

\[ kx = \sin^{-1} y + (kA)y \]  

As figure 2.5 shows, \( y \) becomes double valued near \( kx = (2n + 1)\pi \) when the amplitude of the wave is larger than the wavelength. This is a physically impossible situation and so the wave breaks. Therefore a limiting value of \( E_0 \) is found where \( E_{\text{max}} = n_e e / \epsilon_0 k \). Hence, provided \( \omega \approx \omega_{pe} \) then;

\[ v_E = \frac{\omega_{pe}^2}{k\omega} \approx v_p \]

Therefore the wave breaks if \( v_E \gtrsim v_p \).

The process of wavebreaking generates hot electrons with energies much greater than the Maxwellian average. Electron energies exceeding 10 keV can be generated for very short laser pulses in the order femtoseconds. These are deleterious for many plasma experiments including ICF.
Figure 2.5: Breaking of a wave when its amplitude becomes greater than its wavelength. Here the products $kA = 0$ (solid curve), $kA < 1$ (dashed curve) and $kA > 1$ (dotted curve) are presented.

2.4 Pondermotive force

The momentum associated with an electromagnetic wave can be directly transferred to the plasma due to the pondermotive force. Electrons are more responsive to electric fields than heavy ions. Following Kruer’s discussion, the force equation of a single free electron under the influence of an electromagnetic wave is [43],

$$\frac{\delta v_e}{\delta t} + v_e \cdot \nabla v_e = -\frac{e}{m_e} E(x) \sin(\omega_L t)$$  \hspace{1cm} (2.18)

To the lowest order in $|E|$, $v_e = v^h$ where,

$$\frac{\delta v^h}{\delta t} = -\frac{e}{m_e} E(x) \sin(\omega_L t); \hspace{1cm} v^h = \frac{eE}{m_e \omega_L} \cos(\omega_L t)$$  \hspace{1cm} (2.19)

By averaging the force equation over the rapid electron oscillations we find,

$$m_e \frac{\delta v^s}{\delta t} = -eE^s - m_e \langle v^h \cdot \nabla v^h \rangle_t$$  \hspace{1cm} (2.20)
where the symbol $\langle \rangle_t$ denotes the average over the high frequency oscillations and $v^s = \langle v_e \rangle_t$, $E^s = \langle E \rangle_t$. Substituting for $v^h$,

\[ m_e \frac{\delta v^s}{\delta t} = -eE^2 - \frac{e^2}{m_e \omega L^2} \nabla E^2(x) \]  

(2.21)

The second term in equation 2.21 is called the pondermotive force and is written.

\[ F_p = -\frac{e^2}{4m_e \omega L^2} \nabla E^2 \]  

(2.22)

The pondermotive force has the effect of transporting electrons away from regions of high density to low density. As the force is dependent on $\omega_{pe}$ it is greater for higher electron densities [44]. Consider an electron located at $x_0$. The electron is accelerated due to the force into higher density plasma and stops at $x_1$. In the next cycle it is accelerated back towards the lower density plasma with a greater force, therefore it passes its origin at $x_0$. The next cycle accelerates the electron with a weaker force and therefore it does not reach $x_1$. This action effectively causes a drift in the direction of decreasing wave amplitude, causing a shallower density gradient.

2.5 Self-focussing and filamentation

As a laser pulse travels through an underdense plasma (created by itself, a pre-pulse or otherwise) it will produce a plasma plume centred along the surface normal. There can be a locally higher laser intensity, through a combined effect of thermal forces and the pondermotive force. As a consequence of a lower plasma density the local refractive index in this region is increased and the laser wavefront curves toward the propagation axis. The beam re-focuses, or self-focusses, causing an enhancement of the original density fluctuation. This positive feedback mechanism can lead to the generation of intense filament structures or hot-spots.
which can be detrimental to controlled, uniform irradiation of targets.

Localised high intensity regions may also cause an onset of parametric instability growth [45]. This particular mechanism has been studied numerically, using a non-linear Schrödinger equation, simulated with particle codes, and experimentally verified [46–49]. The threshold for self-focussing and filamentation is given by \((1/k_0\lambda_D)^2(v_{os}/c)^2 > 1\), and is therefore dependent on the intensity of the incident laser beam.

Self-focussing can also occur due to the relativistic increase in electron mass oscillating in an intense electric field. The increase of mass effectively changes the local refractive index. The critical laser power for this to occur is given by [43],

\[
P_{\text{critical}} = 20 \left( \frac{\omega}{\omega_{pe}} \right)^2 \text{[GW]} \tag{2.23}
\]

### 2.6 Inverse Bremsstrahlung

Bremsstrahlung (in German “braking radiation”) is the process where a free electron decelerates in the Coulomb field of an ion and emits a photon. The reverse of this process, where an electron absorbs a photon during a collision with an ion or another electron is known as inverse bremsstrahlung or collisional absorption. At the macroscopic level, electron momentum is driven by an electromagnetic field before being dissipated by means of collisions with ions. Therefore the absorption of laser light by this process is dependent on the electron-ion collision frequency. This is an important process for the transferal of laser energy into the random thermal motion of the plasma. For a plasma in thermal equilibrium Kirchoff’s law can be applied to find the linear (free-free) absorption coefficient \(\alpha_\omega\) [20].

\[
\alpha_\omega(T_e) = \frac{j_\omega(T_e)}{I_{\omega B}(T_e)} \tag{2.24}
\]
Where \( j_\omega \) is the bremsstrahlung emission coefficient and \( I_{\omega B} \) is the spectral intensity at frequency \( \omega \) of a black body at temperature \( T \). This leads to the low intensity expression for the inverse bremsstrahlung absorption rate \( J_{abs} \) [50]:

\[
J_{abs}(J \, s^{-1} \, m^{-3}) = \frac{4}{3} \left( \frac{\pi}{2} \right)^{1/2} \frac{Z^2 n_e n_i e^2 v_{os}^2 \ln \Lambda}{(4\pi\varepsilon_0)^2 [1 - \omega_{pe}^2/\omega_0^2]^{1/2} m_e v_{te}^3} \tag{2.25}
\]

where \( Z \) is the charge state, \( n_i \) is the ion density and \( v_{os} \) the quiver velocity of an electron in an electromagnetic wave. \( \ln \Lambda \) is the Coulomb logarithm, where \( \Lambda \) is the ratio of the maximum and minimum distances of an electron from the ion in a scattering event. The maximum distance is approximately given by the Debye length \( \lambda_D \), whereas the minimum distance is given by the classical distance of close approach or by the de Broglie wavelength of an electron. It can be seen that the absorption by inverse bremsstrahlung is most effective at low temperatures (low \( v_{te} \)) and high densities. The strongest absorption will occur in the densest layer the laser can propagate; at or near the critical surface.

In the previous discussion it was assumed that the radiation field did not perturb the electron velocity distribution significantly. However, in the focus of a high-power laser beam this assumption is not valid. The maximum non-relativistic quiver velocity of an electron in an electromagnetic field can be expressed in terms of the wavelength \( \lambda \) (in \( \mu m \)) and intensity \( I \) (in W cm\(^{-2} \)) of the laser;

\[
v_{os}(cm \, s^{-1}) = \frac{eE_0}{m_e\omega} = 25\lambda I^{1/2} \tag{2.26}
\]

This may exceed the electron thermal velocity. The ratio between the electron quiver velocity and its thermal velocity is referred to as the field strength parameter \( \eta = v_{os}/v_{te} \). If \( \eta \gg 1 \) then the motion of the electrons will be determined primarily by the radiation field. In this case the electron-ion collision time reduces resulting in a lowering of the inverse bremsstrahlung absorption rate. Instead \( J_{abs} \) becomes non-linear, being a function of the irradiance. Many methods have been
used to determine the extent to which inverse bremsstrahlung is reduced due to the increased radiation field. A simple determination is found in [51, 52] where $v_{tc}^2$ is replaced with $v_{te}^2 + v_{os}^2/6$. More recent work suggests that the $\ln \Lambda/v_{te}^2$ term should be replaced with a more complex term [53]. In both cases a very important feature emerges: as the ratio $\eta = v_{os}/v_{te}$ increases to values above unity the power absorbed by inverse bremsstrahlung changes from a linear dependence on the laser intensity $I_L$ to a non-linear $I_L^{-1/2}$ dependence. Therefore this absorption mechanism falls at increased irradiances.

Figures 2.6 and 2.7 show values of $\eta$ and inverse bremsstrahlung heating rates as a function of laser intensity, calculated using a laser wavelength of $\lambda_L = 0.8 \mu m$ with an associated critical density $n_{cr} = 1.7 \times 10^{21} \text{cm}^{-3}$ for electron temperatures representative of a laser produced plasma. Here $Z = 11$, $n_e = n_i = 1 \times 10^{20} \text{cm}^{-3}$, and $\ln \Lambda = 16.3$. It is clear that heating rates follow a quasi-linear relationship for laser intensities up to $\approx 5 \times 10^{14} \text{W cm}^{-2}$ for electron temperatures between 200 eV and 800 eV as the sum of $v_{os}^2 + v_{te}^2 \approx v_{te}^2$. For greater laser intensities there is a sharp decline in inverse bremsstrahlung heating rates. The ratio between the modified and unmodified calculations becomes extremely significant. For example for $I_L = 3 \times 10^{16} \text{W cm}^{-2}$ the ratio for an electron temperature of 200 eV is $\eta \simeq 8$. In this case the linear theory suggests that heating by a 1 ps laser pulse will be 80 eV per electron. As the laser intensity increases collisionless absorption mechanisms, where laser energy is transferred to wave modes in the plasma, becomes the dominant process.

2.7 High frequency electric field amplification

This section introduces the concept of coupling between electromagnetic and plasma waves and the role they have in creating large amplitude plasma waves. Firstly, the process named resonance absorption is described where the compo-
Figure 2.6: The relationship between $\eta$ and laser intensity $I_L$ for a laser wavelength of 0.8 $\mu$m and electron temperatures of 800 eV (blue curve), 400 eV (red curve) and 200 eV (black curve).

Figure 2.7: Inverse bremsstrahlung heating rates for Al$^{+11}$ and laser wavelength of 0.8 $\mu$m. The curves indicate unmodified and modified calculations for electron temperatures of 800 eV (blue curve), 400 eV (red curve) and 200 eV (black curve).
2.7 Laser-Plasma Coupling

Component of the electromagnetic wave into the density gradient causes the electrons in the vicinity of \( n_{cr} \) to oscillate in a resonant manner. Secondly, parametric instabilities are discussed. This mechanism is a decay of a laser photon into two (or more) daughter waves. The daughter wave can be an ion wave, plasmon (plasma wave) or photons depending on parameters such as the density and temperature at which the decay occurs and the electric field strength of the laser. The waves may feedback with one another and become amplified. The process of spectroscopic modulation formation is dependent upon the high frequency electric fields associated with plasmons. Therefore processes such as SBS where the daughter waves are ion acoustic waves and photons are not affecting the Intra-Stark process. We consider two particularly important parametric instability processes where electron plasma waves are involved; SRS and TPD. In this case the laser-plasma coupling is nonlinear in nature and the generated electron plasma waves are exponentially amplified to large values above thermal level plasmons. The electric field associated with thermal level plasmons \( E_{te} \) in units of V cm\(^{-1}\) is,

\[
E_{te} = 3.86 \times 10^{-4} \left( \frac{T_e}{3\pi k} \right)^{1/2}
\]

(2.27)

Where the electron temperature \( T_e \) is expressed in eV and the plasmon wave number is expressed in cm\(^{-1}\) [20].

2.7.1 Resonance absorption

Electron plasma waves can be created directly via the absorption of electromagnetic radiation in the plasma medium. An example of this occurs when a \( p \)-polarized light wave incident into an inhomogeneous plasma is reflected at the critical surface, with a turning point determined by \( n_e = n_{cr} \cos^2 \theta \) where \( \theta \) is the angle of incidence of the wave. The electric field of the wave at the turning point is in the same direction of the plasma density gradient \( \nabla n_e \). A portion of the field
tunnels through the critical surface region and resonantly drives the electrons to oscillate at $\omega_{pe}$ causing heating of the plasma, a process called resonance absorption. This tends toward a maximum value for angles $\theta \simeq 60^\circ$ which correspond to densities $n_e \simeq n_{cr}/4$. The resonantly generated plasma waves travel away from regions of high electron density ($n_e > n_{cr}$) towards regions of lower density. The electrostatic field becomes sufficiently intense and localized so that wavebreaking occurs. Here, a portion of the electrons at the critical surface are accelerated through the localized field in less than an oscillation period. The fraction of high energy electrons, or hot electrons, to thermal electrons is dependent on the product $I\lambda^2$ and can be significant (in excess of 20%) for intensities above $10^{18}$ W cm$^{-2}$ [20].

In the resonant layer the field grows linearly in time and becomes localised. The driving electric field $E_D$ in this resonant layer is given by;

$$\frac{E_D^2}{E_0^2} = 0.44 \frac{f_R}{k_0L}$$  \hspace{1cm} (2.28)

where $E_0$ is the amplitude of the oscillating laser electric field, $E_0 \cos(\omega t)$ [44];

$$E_0 (\text{V cm}^{-1}) = 2.75 \times 10^9 \left( \frac{I_L}{10^{16}\text{W cm}^{-2}} \right)^{1/2}$$  \hspace{1cm} (2.29)

and $f_R$ is the total approximated resonant absorption coefficient for a cold plasma [54]. $f_R$ is dependent on the density scale length $L_n$, wavenumber of the laser light $k_0$ and the angle of propagation into the plasma. Forslund et al discovered that in a hot plasma $f_R$ is almost independent of the electron temperature up to values of $T_e \sim 100$ keV, with a maximum value of 50 % [55]. The presence of large amplitude electron plasma waves when resonant with the laser produces a pondermotive force which expels plasma from the region and steepens the local density gradient and may generate a rippled critical density surface.

As the wave grows in time it will reach a point where it will either break,
2.7 Laser-Plasma Coupling

Figure 2.8: Schematic of resonance absorption

as discussed in section 2.3.2, or the process of convection will halt any further growth. Wave breaking and convection will limit the value of the electric field to $E_B$ or $E_C$ respectively;

$$E_B = \frac{m_e \omega_0}{e} \left( \frac{2eE_D}{m_e \omega_0} \omega_0 L \right)^{1/2}$$
(2.30)

$$E_C = 1.2E_D \left( \frac{\omega_0 L}{v_{te}} \right)^{2/3}$$
(2.31)

2.7.2 Parametric instabilities

Parametric instabilities can be represented as the resonant coupling of the incident electromagnetic wave, or pump wave, into two high frequency daughter waves; electron plasma waves and scattered electromagnetic waves. Consider two existing modes of oscillation within the plasma with frequencies $\omega_1$ and $\omega_2$ and wave numbers $k_1$ and $k_2$ which have amplitudes at thermal level (commonly called noise level). The pump wave with a frequency $\omega_0$ and wavenumber $k_0$ will beat with the two thermal waves producing side-bands at $\omega_1 = \omega_0 - \omega_2$ and at $\omega_2 = \omega_0 - \omega_1$. Therefore both frequencies $\omega_1$ and $\omega_2$ will re-enforce one another as
they beat with increased strength with the pump wave. The amplified modes will grow as they continue to receive energy from the pump wave. Losses are due to Landau damping. Therefore a threshold must be exceeded for the growth of the instabilities to occur. If the instability grows in time at a single localized point then it is said to be an absolute instability. In this case the instability grows without the need of an external feedback mechanism (reflections for example) and instead grows due to internal feedback. Convective instabilities, on the other hand, grow as they propagate along the system and require external feedback to be able to grow exponentially in time. The disturbance will eventually dissipate if observing a fixed point and therefore it will obey dispersion relations dependent on the matching conditions. The matching conditions require that energy and momentum are conserved;

\[ \omega_0 = \omega_1 + \omega_2 \quad \text{and} \quad k_0 = k_1 + k_2 \quad (2.32) \]

### 2.7.3 Stimulated Raman Scattering

Stimulated Raman Scattering (SRS) is the resonant decay of an electromagnetic wave into a scattered electromagnetic wave and an electron plasma wave, denoted by the subscripts \(sc\) and \(epw\) respectively. From the matching conditions the minimum frequency of the electromagnetic wave is the electron plasma frequency of the local electron density. Therefore it is found that SRS takes place at electron densities \(n_e \leq n_{cr}/4\). Therefore, when SRS occurs, a broadband frequency of electron plasma waves are amplified. The frequency width is dependent on the electron density gradient and the extent of the plasma. Detection of the scattered SRS electromagnetic wave gives information regarding the generation of the instability itself and the density of the plasma. The matching conditions are;

\[ \omega_0 = \omega_{sc} + \omega_{epw} \quad \text{and} \quad k_0 = k_{sc} + k_{epw}\quad (2.33) \]
This process involves the deposition of part of the incident laser light into an electron plasma wave, the remainder energy is scattered. The portion of energy absorbed into the electron plasma wave will heat the plasma as the electron wave damps. Consider light waves with electric field amplitude \( E_L \) propagating through a plasma density fluctuation associated with an electron plasma wave. In this case the electrons oscillate due to the presence of the laser light and a transverse current is created. For instances where the laser and electron plasma wave frequencies and wavenumbers match, the transverse current generates a scattered light wave. This scattered light interferes with the incident laser light and there is a change in the wave pressure. This generates a feedback loop where a small density fluctuation creates a transverse current, generating a scattered light wave and reinforcing the density fluctuation.

To examine the characteristics of SRS, Kruer’s work is followed closely [43]. Beginning with Ampere’s law,

\[
\nabla \times \mathbf{B} = \frac{4\pi}{c} \mathbf{J} + \frac{1}{c} \frac{\delta \mathbf{E}}{\delta t} \tag{2.34}
\]

The electric and magnetic terms can be expressed in terms of the vector potential \( \mathbf{A} \) and electrostatic potential \( \phi \), where \( \mathbf{B} = \nabla \times \mathbf{A} \). Substituting for \( \mathbf{B} \) and \( \mathbf{E} \) and choosing \( \nabla \cdot \mathbf{A} = 0 \) then,

\[
\left( \frac{1}{c^2} \frac{\delta^2}{\delta t^2} - \nabla^2 \right) \mathbf{A} = \frac{4\pi}{c} \mathbf{J} - \frac{1}{c} \frac{\delta}{\delta t} \nabla \phi \tag{2.35}
\]

The current density \( \mathbf{J} \) can be separated into a transverse part \( \mathbf{J}_t \) associated with the laser light and a longitudinal part associated with the electron plasma wave \( \mathbf{J}_l \). Poisson’s equation and the equation of charge can be used to relate \( \mathbf{J}_l \) to \( \nabla \phi \),

\[
\nabla^2 \phi = -4\pi \rho \tag{2.36}
\]

\[
\frac{\delta \rho}{\delta t} + \nabla \cdot \mathbf{J} = 0 \tag{2.37}
\]
Since $\nabla \cdot \mathbf{J}_t = 0$ then,

$$\frac{\delta}{\delta t} \nabla \phi = 4\pi \mathbf{J}_t \tag{2.38}$$

Thus equation 2.35 becomes,

$$\left( \frac{1}{c^2} \frac{\delta^2}{\delta t^2} - \nabla^2 \right) \mathbf{A} = \frac{4\pi}{c} \mathbf{J}_t \tag{2.39}$$

The transverse current can be expressed as $\mathbf{J}_t = -n_e \mathbf{v}_{os}$ where $\mathbf{v}_{os}$ is the oscillation velocity of an electron in the electric field of a light wave. As $\mathbf{v}_{os} << c$ then $\mathbf{v}_{os} = e\mathbf{A}/m_ec$.

The scattering of a large amplitude light wave $\mathbf{A}_L$, such as laser light, by a small amplitude density fluctuation $\tilde{n}_e$ associated with an electron plasma wave can be determined by considering the scattered amplitude $\mathbf{A} = \mathbf{A}_L + \tilde{\mathbf{A}}$ and the new density due to the perturbation $n_{scat} = n_0 + \tilde{n}_e$ where $n_0$ is the uniform background plasma density. Using 2.39,

$$\left( \frac{\delta^2}{\delta t^2} - \frac{c^2}{\omega_{pe}^2} \nabla^2 \right) \tilde{\mathbf{A}} = -\frac{4\pi e^2}{m_e} \tilde{n}_e \mathbf{A}_L \tag{2.40}$$

To obtain the equation for the density fluctuation associated with an electron plasma wave ion motion is neglected due to their large mass. The continuity and force equations are then,

$$\frac{\delta n_e}{\delta t} + \nabla \cdot (n_e \mathbf{v}_{os}) = 0 \tag{2.41}$$

$$\frac{\delta \mathbf{v}_{os}}{\delta t} + \mathbf{v}_{os} \cdot \nabla \mathbf{v}_{os} = -\frac{e}{m_e} \left( \mathbf{E} + \frac{\mathbf{v}_{os} \times \mathbf{B}}{c} \right) - \frac{\nabla p_e}{n_em_e} \tag{2.42}$$

where $p_e$ is the pressure of an electron fluid. Using $\mathbf{v}_{os} = \mathbf{v}_L + e\mathbf{A}/m_ec$ to separate the longitudinal and transverse components equation 2.40 becomes,

$$\frac{\delta \mathbf{v}_L}{\delta t} = \frac{e}{m} \nabla \phi - \frac{1}{2} \nabla \left( \mathbf{v}_L + \frac{e\mathbf{A}}{m_ec} \right)^2 - \frac{\nabla p_e}{n_em_e} \tag{2.43}$$
To generate an equation which describes the generation of fluctuations in the electron density by electromagnetic intensity variations the values $v_L = \tilde{v}$, $n_e = n_0 + \tilde{n}_e$, $A = A_L + \tilde{A}$ and $\phi = \tilde{\phi}$ are used. The tilde denotes an infinitesimal quantity. Using the adiabatic equation of state and linearising 2.41 and 2.42,

\[ \frac{\delta \tilde{n}_e}{\delta t} + n_0 \nabla \cdot \tilde{\mathbf{v}} = 0 \]  
\[ \frac{\delta \tilde{\mathbf{v}}}{\delta t} = \frac{e}{m_e} \nabla \tilde{\phi} - \frac{e^2}{m_e^2 c^2} \nabla (A_L \cdot \tilde{A}) - \frac{3v_{te}^2}{n_0} \nabla \tilde{n}_e \]

Where $v_{te}$ is the electron thermal velocity. Thus the electron density fluctuation becomes,

\[ \left( \frac{\delta^2}{\delta t^2} + \omega_{pe}^2 - 3v_{te}^2 \nabla^2 \right) \tilde{n}_e = \frac{n_0 e^2}{m_e^2 c^2} \nabla^2 (A_L \cdot \tilde{A}) \]  

The maximum growth rate of the SRS instability can be calculated by considering the dispersion relation. Here it is necessary to use $A_L = A_0 \cos(k_0 \cdot \mathbf{x} - \omega_0 t)$ and to Fourier analyse equations 2.40 and 2.46.

\[ (\omega^2 - k^2 c^2 - \omega_{pe}^2) \tilde{\mathbf{A}}(k, \omega) = \]

\[ \frac{4\pi e^2}{2m_e} A_0 [\tilde{n}_e(k - k_0, \omega - \omega_0) + \tilde{n}_e(k + k_0, \omega + \omega_0)] \]

\[ (\omega^2 - \omega_{ek}^2) \tilde{n}_e(k, \omega) = \]

\[ \frac{k^2 e^2 n_0}{2m_e^2 c^2} A_0 \cdot [\tilde{\mathbf{A}}(k - k_0, \omega - \omega_0) + \tilde{\mathbf{A}}(k + k_0, \omega + \omega_0)] \]

Where $\omega_{ek}$ is the Bohm-Gross frequency (equation 2.5). Eliminating $\tilde{\mathbf{A}}$ in equation 2.48 using 2.47, equating $D(\omega, k) = \omega^2 - k^2 c^2 - \omega_{pe}^2$ and neglecting the terms $\tilde{n}_e(k - 2k_0, \omega - 2\omega_0)$ and $\tilde{n}_e(k + 2k_0, \omega + 2\omega_0)$ as nonresonant, the dispersion equation for SRS becomes,

\[ \omega^2 - \omega_{ek}^2 = \frac{\omega_{pe}^2 k^2 v_{te}^2}{4} \left[ \frac{1}{D(\omega - \omega_0, k - k_0)} + \frac{1}{D(\omega + \omega_0, k - k_0)} \right] \]  

45
using $\omega = \omega_{ek} + i\gamma$ the growth rate of the instability becomes,

$$\gamma = \frac{k\nu_{os}}{4} \left[ \frac{\omega_{pe}^2}{\omega_{ek}(\omega_0 - \omega_{ek})} \right]^{1/2}$$

(2.50)

For $n_e << n_{cr}/4$ the wavenumber becomes $k = 2k_0$, and for $n_e \sim n_{cr}/4$ the wavenumber becomes $k = k_0$.

The threshold condition which must be met for SRS to exist can be found by considering the density scale length through which the laser light must propagate. Propagation of light away from the matching region introduces an effective damping rate of approximately $v_{g1}v_{g2}/K'$ where $v_{g1}$ and $v_{g2}$ are the group velocity of the growing waves and $K'$ is the spatial derivative of the wavenumber mismatch. Therefore,

$$\left| \frac{\gamma^2}{K'v_{g1}v_{g2}} \right| \geq 1$$

(2.51)

Neglecting temperature gradients and assuming a linear variation on density scale length and noting $v_{g2} = c$ then the threshold for SRS becomes,

$$\left( \frac{\nu_{os}}{c} \right)^2 > \frac{2}{k_0L}$$

(2.52)

### 2.7.4 Two Plasmon Decay

Two Plasmon Decay (TPD) is the process where the incident pump wave decays into two electron plasma waves, denoted by the subscripts $epw1$ and $epw2$. As both the daughter electron plasma waves are approximately equal to the local plasma frequency $\omega_{pe}$, the instability takes place in a localised region at densities $n_e \sim n_{cr}/4$. Therefore the wave amplification will have a narrow frequency spectrum. Large amounts of hot electrons are produced from this interaction by the subsequent damping of the electron plasma waves. Applying the matching
conditions we have;

\[ \omega_0 = \omega_{epw1} + \omega_{epw2} \quad \text{and} \quad k_0 = k_{epw1} + k_{epw2} \]  \hspace{1cm} (2.53)

To derive the instability ion motion is neglected due to their large mass. Using the same notation as the previous subsection and following Kruer’s work then [43],

\[ \frac{\delta \tilde{n}_e}{\delta t} + n_0 \nabla \cdot \tilde{v}_L + \mathbf{V}_{os} \cdot \nabla \tilde{n}_e = 0 \]  \hspace{1cm} (2.54)

\[ \frac{\delta \tilde{v}_L}{\delta t} = \frac{e}{m_e} \nabla \tilde{\phi} - \frac{3v_{te}^2}{n_0} \nabla \tilde{n}_e - \nabla (\mathbf{v}_{os} \cdot \tilde{v}_L) \]  \hspace{1cm} (2.55)

The derivative with respect to time of equation 2.54 and the divergence of equation 2.55 gives,

\[ \frac{\delta^2 \tilde{n}_e}{\delta t^2} + (\omega_{pe}^2 - 3v_{te}^2 \nabla^2)\tilde{n}_e + \frac{\delta (\mathbf{v}_{os} \cdot \nabla \tilde{n}_e)}{\delta t} - n_0 \nabla^2 (\mathbf{v}_{os} \cdot \tilde{v}_L) = 0 \]  \hspace{1cm} (2.56)

Representing \(\mathbf{v}_{os} = v_{os} [\exp(i\mathbf{k}_0 \cdot \mathbf{x} - i\omega_0 t) + \exp(-i\mathbf{k}_0 \cdot \mathbf{x} + i\omega t)]/2\) and Fourier analysing 2.56,

\[ (-\omega^2 - \omega_{ek}^2)\tilde{n}_e(k, \omega) \]

\[ + \frac{\omega}{2} \mathbf{k} \cdot \mathbf{v}_{os}[\tilde{n}_e(k - k_0, \omega - \omega_0) + \tilde{n}_e(k + k_0, \omega + \omega_0)] \]

\[ + \frac{n_0 k^2}{2} \mathbf{v}_{os} \cdot [\tilde{v}_L(k - k_0, \omega - \omega_0) + \tilde{v}_L(k + k_0, \omega + \omega_0)] = 0 \]

Choosing \(\omega = \omega_{pe}\), neglecting any off-resonant terms and knowing that the the approximate value of \(\tilde{v}_L\) is,

\[ \tilde{v}_L \simeq \frac{k}{k^2 \omega} \frac{\tilde{n}_e(k, \omega)}{n_0} \]  \hspace{1cm} (2.58)

The dispersion relation for TPD can now be stated as,

\[ (\omega^2 - \omega_{ek}^2)(\omega - \omega_0)^2 - \omega_{ek-k_0}^2) = \left[ \frac{k \cdot \mathbf{v}_{os} \omega_{pe} [(k - k_0)^2 - k^2]}{2k | \mathbf{k} - \mathbf{k}_0 |} \right]^2 \]  \hspace{1cm} (2.59)
The maximum growth rate is found by substituting $\omega = \omega_{ek} + i\gamma$, analysing the frequency matching and setting $k >> k_0$ for electron plasma waves propagating at $45^\circ$ to both $k_0$ and $v_{os}$:

$$\gamma \simeq \frac{k_0v_{os}}{4}$$

(2.60)

The threshold for TPD is found by taking into consideration the plasma inhomogeneity;

$$\left(\frac{v_{os}}{v_{te}}\right)^2 \simeq \frac{12}{k_0L}$$

(2.61)

Note that the threshold for TPD is much lower than that associated with SRS.

### 2.7.5 Convective amplification

Using the work of Sinha and Gupta, the interaction length $L_{int}$ over which the amplification takes place is limited by the wavenumber mismatch $K(x)$ and is expressed as $L_{int} = (|K'|)^{-1/2}$ [19]. As the plasmons propagate the interaction length, their amplitude increases by a factor

$$A = A_{te} \exp A_C$$

(2.62)

where $A_{te}$ is the thermal level amplitude of plasmons and $A_C$ is the maximum convective amplification coefficient in the interaction region. $A_C$ is written as,

$$A_C = \frac{\pi\gamma_0^2}{|K'v_{1x}v_{2x}|}(1 - \Gamma_C/\gamma_0)$$

(2.63)

Here $\gamma_0$ is the growth rate, $\Gamma_C$ is the collisional damping rate of plasmons and $v_{1x}$ and $v_{2x}$ are the $x$ components of the group velocities of the plasmons. If Landau damping is insignificant for example where $k\lambda_D < 0.3$ then the collisional damping rate is dominant. This can be expressed as $\Gamma_C = \nu_{ei}/2$ where $\nu_{ei}$ is the
2.8 Laser-Plasma Coupling

Electron-ion collision frequency expressed as [19],

\[ \nu_{ei} = 6.5 \times 10^{-6} \bar{Z} n_e T_e^{-3/2} \ln \Lambda \] (2.64)

\[ \Lambda = \frac{9.92 \times 10^5 T_e^{5/2}}{\bar{Z}^{3/2} n_e^{1/2}} \] (2.65)

where \( \bar{Z} \) is the average ionisation of an ion and \( \ln \Lambda \) is the Coulomb logarithm.

The maximum convective amplification coefficient in an inhomogeneous plasma is

\[ A_C = \frac{\pi}{6} \frac{|v_{os}|^2}{v_{te}^2} (k_0 L_n) \left( 1 - \frac{2 \Gamma_C}{k_0 |v_{os}|} \right) \] (2.66)

The presence of instability growth is detrimental for inertial fusion applications where the interaction of laser light with preformed plasma with large density gradients occurs. The production of high phase velocity plasma waves can lead to the creation of high energy hot electrons and preheating of the fusion fuel. The fuel must be on a low adiabat to ensure efficient compression. For electromagnetic daughter waves the laser energy will scatter from the target reducing the energy absorption in the plasma. Furthermore, as the majority of instabilities occur at densities significantly less than the critical density the laser energy will couple into the underdense region. Overall these characteristics will reduce the fuel compression and fusion gain. Experimentally it has been found that a large fraction of the laser energy (up to 50\%) may be coupled into parametric instabilities [20].

2.8 Magnetic field generation

The presence of self-generated strong magnetic fields in plasma can affect both hydrodynamic behaviour, in terms of the transport of hot electrons, energy and ions [56], and also the emission from bound-bound transitions where the Zeeman effect causes the splitting of atomic levels [57, 58]. The generation of magnetic
fields in plasma is due to the non-vanishing curl of the electric field $E$. The volume forces due to the free electrical charges balance any gradient in the electron pressure, $0 = -\nabla P_e - e n_e E$. Using the Maxwell equation and the ideal gas equation to find the pressure term;

$$\frac{\partial B}{\partial t} = \frac{c k_B}{e n_e} \nabla T_e \times \nabla n_e \tag{2.67}$$

The magnitude of the electric field $B$ can be found by making the substitution;

$$\frac{\partial}{\partial t} \approx \frac{1}{\tau_L} \quad \frac{\nabla n_e}{n_e} \approx \frac{1}{L_n} \quad \nabla T_e \approx \frac{T_e}{L_T} \tag{2.68}$$

into equation 2.67, giving;

$$B (\text{MGauss}) \approx 10 \left( \frac{\tau_L}{1 \text{ ns}} \right) \left( \frac{k_B T_e}{1 \text{ keV}} \right) \left( \frac{30 \mu m}{L_n} \right) \left( \frac{30 \mu m}{L_T} \right) \tag{2.69}$$

where $\tau_L$ is the laser pulse duration or the loss time due to ablation and diffusion, $L_n$ is the density gradient scale length and similarly $L_T$ is the temperature.
2.9 Laser-Plasma Coupling

Therefore a magnetic field is created due to $\nabla n_e \times \nabla T_e$, i.e. non-parallel density and temperature gradients. Similarly, the role of return currents due to convective plasma flow produces extra terms into equation 2.69. By considering the plasma as a magnetic fluid it is possible to develop a generalised Ohm’s law with a characteristic rate of change of magnetic field. Two terms are found to dominate; the $\nabla n_e \times \nabla T_e$ term, and the convective plasma flow term $\nabla \times (\mathbf{u} \times \mathbf{B})$ where $\mathbf{u}$ is the centre of mass fluid flow. Weaker contributions include the $\mathbf{J} \times \mathbf{B}$ Hall effect and the voltage drop $\mathbf{J}/\sigma_E$ where $\sigma_E$ is the plasma electrical conductivity. Considering a constant current solution ($\partial B/\partial t$) and letting $\nabla \approx 1/L_P$ and $u \approx c_s$, where $L_P$ is the pressure gradient scale length and $c_s$ is the sound speed, and using the ideal gas equation of state we find:

$$B(\text{MGauss}) \approx \left( \frac{m_i}{m_p} \right)^{1/2} \left( \frac{T_e}{\text{keV}} \right)^{1/2} \left( \frac{30 \mu m}{L_P} \right)$$

(2.70)

where $m_i$ and $m_p$ are the ion and proton masses respectively. This implies that for an electron temperature of 1 keV and a pressure scale length of 30 µm toroidal magnetic field of 1 MGauss is generated within the plasma.

2.9 Summary

In this chapter the spectrum of waves that can be supported by a plasma was presented including electromagnetic waves, electron plasma waves (plasmons) and ion acoustic waves. The high frequency electric fields associated with the plasmons is important in Intra-Stark formation. Dispersion relations were derived using a fluid and kinetic treatment. The maximum amplitude of plasmons by wave breaking and particle trapping was described. Amplification of plasmon modes by non-linear laser-plasma coupling mechanisms was given. The important processes of TPD and SRS are identified. These mechanisms can occur for laser intensities as low as $5 \times 10^{13}$ W cm$^{-2}$. Wave damping by Landau and col-
lisional processes was discussed. Filamentation of laser light within the plasma was presented where hot spots are created due to refraction. This is dependent on the intensity of the incident laser beam. Laser induced magnetic fields were presented and examples given.
Chapter 3

Plasma spectroscopy

3.1 Introduction

Plasma spectroscopy is a powerful tool for determining the parameters and characteristics of plasma. It is impossible to probe high density plasma with laser light where \( n_e > n_{cr} \) for a given laser wavelength. This is typical in laser plasma interactions. Information such as electron density, temperature and electric fields can be measured by observing and analysing the X-ray emission from high density plasma. The lineshape of the dispersed emitted radiation also contains information about high frequency waves supported by the plasma. The interaction between waves and the radiators creates structures on the emitted lineshape which are a signature of the wave amplitudes and frequencies. The wave amplitudes and frequencies provide information about the laser-plasma coupling mechanisms from which they were created.

This chapter describes the processes by which emission from atomic radiators is affected by the fields in the plasma medium. Firstly radiative processes and the extent to which they are affected by pressure and Stark broadening are outlined. This includes thermal ion motion and the distribution of quasistatic
ion-microfields. The effect of strong fields is then discussed with special reference to the high frequency electric fields associated with plasmons. The interaction between the low frequency quastistatic ion-microfields and the high frequency plasmons is described. This field is called “Intra-Stark Spectroscopy” [30, 31]. The effects this interaction upon the line shape is presented in the case of a high density plasma.

3.2 Radiative processes

The ionisation state and radiation field of a plasma is determined by the action of a strong electric field associated with laser light upon it and the subsequent collisional, recombinational and radiative interactions between its constituent particles, ions and electrons. Elastic and inelastic collisions result in the imparting of energy from one particle to another. Elastic collisions lead to an equilibration of particle kinetic energies due to their conservation during the collision. The distribution of the particle velocities follows a Maxwellian form characterised by the temperature of the particle species $T_s$. During inelastic collisions energy is transferred from the colliding particle to the bound states of the ion. Alongside the action of laser light, including above threshold ionisation and multi-photon ionisation, inelastic collisions result in the movement of electron populations in ground and excited bound and free states [59, 60]. The principal mechanisms of this process are listed in table 3.1.

As a result of electrons being promoted into excited states and subsequently depopulating into lower energy states a radiation field is generated within the plasma. This consists of transitions between free-free (bremsstrahlung radiation), free-bound (recombination radiation) and bound-bound states (line radiation). The study of line radiation reveals a great deal of information about the plasma. The presence and relative intensity of line radiation from a plasma is sensitive
### 3.3 Plasma spectroscopy

<table>
<thead>
<tr>
<th>Process</th>
<th>Equation</th>
</tr>
</thead>
<tbody>
<tr>
<td>Collisional ionisation</td>
<td>$(n_i^z + e \rightarrow n_g^{z-1} + e + e)$</td>
</tr>
<tr>
<td>Three-body recombination</td>
<td>$(n_g^{z+1} + e + e \rightarrow n_j^z + e)$</td>
</tr>
<tr>
<td>Collisional excitation</td>
<td>$(n_i^z + e \rightarrow n_j^z + e)$</td>
</tr>
<tr>
<td>Collisional de-excitation</td>
<td>$(n_j^z + e \rightarrow n_i^z + e)$</td>
</tr>
<tr>
<td>Dielectronic recombination</td>
<td>$(n^z + e \rightarrow n_{j,j'}^{z-1} \rightarrow n_j^{z-1} + \hbar\omega)$</td>
</tr>
</tbody>
</table>

Table 3.1: Inelastic collisional processes between particles. Here $n$ denotes the state of the ion, the superscript $z$ is the number of electrons within the bound states, subscript $i$, $j$ and $g$ are the lower and upper excited states and ground state respectively. During dielectronic recombination $n_{j,j'}^{z-1}$ denotes an ion in a doubly excited state.

to the ionisation stage, the ion number density and the population of electrons within the excited state under observation. The width and structure of the line itself is an indication of the density and temperature of the plasma from which it has escaped. Strong electric or magnetic fields generated within the plasma can perturb the transition between bound states by lifting the degeneracy and/or splitting the energy levels, depending on the configuration of the electrons in their states and the lifetime of the fields.

### 3.3 Spectral line broadening

The photon energies related to the transitions between bound states of stationary, isolated atoms or ions have an energy “spread”. This is determined by the Einstein A-rates and the limits of the uncertainty principle, i.e. the natural line width ($\hbar \approx \Delta E\Delta t$). Within a plasma environment in thermal equilibrium ions have a characteristic speed dependent on the local ion temperature $T_i$. Therefore in the experiment (stationary) frame of reference the ions emit bound-bound photon energies that are Doppler shifted by an amount dependent on the temperature of the ions and line of sight. Assuming the radiator (ion) velocities are non-relativistic, which they commonly are due to the large mass of the ion, they will follow a Maxwellian distribution. The corresponding line shape function, a parameter that determines the formation of the line accounting for perturbations
on the radiators, is Gaussian,

\[ L_D(\omega) = (\sqrt{\pi} \omega_D)^{-1} \exp \left[ - \left( \frac{\Delta \omega}{\omega_D} \right)^2 \right] \]  

(3.1)

and the Doppler broadening parameter, the half width at half maximum (HWHM) of the Gaussian line shape, is given by,

\[ \omega_D = \left( \frac{2k_B T_i}{m_i c^2} \right)^{1/2} \omega_s \]  

(3.2)

where \( \Delta \omega \) is the frequency detuning from the stationary frame transition frequency \( \omega_s \). The broadening effects due to the natural line width is much less than Doppler broadening and so it can be neglected in the derivation.

Pressure broadening is the action of other particles within the plasma perturbing the radiator. In a plasma the primary cause of pressure broadening is due the Stark effect on the bound state energy levels of radiators caused by electric fields. These fields can be generated by neighbouring electrons or ions, or by the collective fields associated with plasma waves. The mechanisms can be split into two categories, the Holtsmark approximation for heavy, slow particles (quasistatic) and Lorentz approximation for light, fast particles (impact). These fields must have a lifetime greater than the upper state lifetime upon which they act.

The collective fields generated by the charged ions are called ion-microfields. The Holtsmark approximation describes the quasistatic ion-microfields \( F \) within the plasma interacting with the radiators. Holtsmark considered the distribution of the field strength magnitudes \( W(F) \) as isotropic. The Coulomb fields are considered to be unshielded with a uniform distribution. Then the interaction between neighbouring ions causes a reduction in the field strength. This is named the reduced field strength. The distribution of field strength magnitudes can then
be written as [61],

$$W_H(F) = H_0(\beta)/F_0$$  \hspace{1cm} (3.3)

where $\beta = F/F_0$ is the reduced field strength, $H_0(\beta)$ is the Holtsmark function and $F_0$ is the Holtsmark field strength,

$$F_0 = 2.603 \frac{Ze}{4\pi\epsilon_0} n_i^{2/3}$$  \hspace{1cm} (3.4)

Therefore the ion-microfield distribution is dependent on the mean ion-ion separation, the charge of the ions and the Debye shielding length $\lambda_D$ which governs the extent to which the Coulomb field is shielded. The ion-microfield distribution function $H_0(\beta)$ is shown in figure 3.1 as a function of the reduced field strength [62].

As the time scale for changes in the Holtsmark field is usually much longer than the transition time for allowed transitions between bound states it manifests itself as a Stark broadening of the bound states and, therefore, the emitted line profile. For example the radiative rate of the Al$^{+11}$ ($1s3p \rightarrow 1s^2$) He $\beta$ transition is $7.72 \times 10^{12}$ s$^{-1}$ whereas the lifetime of the Holtsmark field has a time scale of the order of the hydrodynamic plasma expansion. The ion-microfield distribution function is quasistatic. There is no simple description of line shapes. The field
Figure 3.2: Spectral line shapes are presented for Doppler (Gauss) and impact (Lorentz) broadening, and as a convolution (Voigt).

Lorentz’s approximation describes the relation between the full width half maximum (FWHM) of spectral lines and the rate of effective electron collisions. The process of an electron colliding with a radiator interrupts the emission process and broadens the line profile of the transition. This occurs provided the electron collision duration is short compared to the decay time of the radiator. This is named the Impact Approximation. The corresponding line shape function created by impact broadening is Lorentzian and is expressed,

$$L_L(\omega) = \frac{w/\pi}{w^2 + (\Delta\omega - d)^2}$$  \hspace{1cm} (3.5)

where $w$ is the half width half maximum (HWHM) of the line under consideration and $d$ is the shift of the line centre from $\omega_c$. In some cases where Doppler and impact broadening are equally important it is not possible to treat the two processes independently when calculating the overall line shape function. The combined action can be expressed as a convolution, which is in this case a Voigt profile (see figure 3.2). In effect the impact broadening effects the wings of the profile and the quasistatic broadening effects the line centres.
3.4 Effects of strong fields on radiative transitions

This section will describe how the strong magnetic and electric fields generated within plasmas (or externally applied) can alter the characteristic energy levels of the bound states. This includes both the Zeeman splitting of lines and the quadratic and linear Stark effect. Both effects may be observed as features (peaks or dips) near or on the resonant line under investigation. These phenomena offer insights into how the plasma evolves hydrodynamically and can reveal the coupling effects between laser and plasma.

3.4.1 The Zeeman effect

The process by which the atomic energy levels are split into separate components in a magnetic field is the Zeeman effect. A magnetic field completely removes the degeneracy of the levels with respect to the projections $M_J$ of the total angular momentum $J$. The magnetic moment of the atom $\mu = -\mu_B g_J J/\hbar$, where $\mu_B$ is the Bohr magneton and $g_J$ is the gyromagnetic ratio or Landé factor, is acted upon by a magnetic field splitting the $J$ level into $2J+1$ symmetrical components, $M_J = 0, \pm1, \pm2, \ldots \pm J$. The magnitude of the splitting is dependent on the magnetic field strength, the projection quantum number $M_J$ under consideration and the Landé factor of the atomic level,

$$g_J = 1 + \frac{J(J+1) - L(L+1) + S(S+1)}{2J(J+1)}$$

where $L$ and $S$ denote the total orbital and spin angular momentum quantum number of an atom. Using the $LS$ coupling approximation the spectral lines emitted from a magnetically perturbed radiator contain extra features. This approximation can be used for highly ionised systems. The extra features are
created by new transition channels created by the splitting, the spacing between one another dependent of the field strength. Transitions from an excited state to a less excited state will result in a more complex spectra due to the splitting of both the upper and lower levels. The Al\(^{+11}\) \((1s3p \rightarrow 1s^2)\) He\(\beta\) transition is shown in figure 3.3 as a function of the magnetic field imposed upon the radiator. Here the upper state (singlet) is split into the components \(M_J = 0, \pm 1\) while the lower state is left uneffected.

### 3.4.2 The Stark effect

The Stark effect describes the action of an external electric field upon energy levels of an atom/ion. For helium like ions where the atom is stripped of all but two electrons the Stark effect can be split into two categories. These are the quadratic Stark effect for weak electric fields and linear Stark effect for strong electric fields. For hydrogen like ions the linear Stark effect applies for weak electric fields due to the degeneracy of the angular momentum quantum number \(l\). Because of this the atomic states in hydrogenic ions will be more susceptible to the electric field [63]. In these instances the linear Stark effect occurs at low electric field strengths with respect to atoms/ions with more than one bound.
electron. In multi-electron ions an increase in field strength (calculated using the quadratic shift) becomes comparable with the energy splitting between terms. The mixing of states with different parities becomes significant and the quadratic Stark effect becomes linear. In other words the linear Stark effect will take place when fine structure splitting due to the bound electron-electron interactions is less than the splitting caused by the external field. Therefore for a given excited state the external electric field must exceed a threshold value for this to occur.

The perturbation in an energy level due to the quadratic Stark effect is,

$$E^{(2)}_{nlm_l} = F_{ex}^2 \chi \frac{9n^2}{4(Z-1)^2(2l+1)}$$

(3.7)

where,

$$\chi = \left[ \frac{(n^2-(l+1)^2)((l+1)^2-m_l^2)}{(2l+3)(E_{nl}-E_{nl+1})} + \frac{(n^2-l^2)(l^2-m_l^2)}{(2l-1)(E_{nl}-E_{nl-1})} \right]$$

(3.8)

Here $n$, $l$, and $m_l$ are the principal, angular and magnetic quantum numbers of the electron respectively, $F_{ex}$ is the external electric field, and $Z$ is the nuclear charge of the ion. In this approximation there is no Stark effect upon the ground state. It is clear that the maximum Stark shift will occur for for $m_l = 0$.

Figure 3.4 presents an example of the splitting induced on the $\text{Al}^{+11} \, ^1S$, $^1P$ and $^1D$ $n = 3$, $m_l = 0$ states. The $^1S$ level remains relatively unchanged at high field strengths due to the absence of angular momentum. In contrast the $^1P$ and the $^1D$ levels begin to separate energetically from one another at high field strengths. This occurs for two reasons; the $^1P$ level has a higher unperturbed energy state than $^1D$, and the second term becomes dominant in equation 3.8 for the $^1D$ level.

The ratio between the $^1P$ and $^1D$ $n = 3$, $m_l = 0$ states for $\text{Mg}^{+10}$, $\text{Al}^{+11}$ and $\text{Si}^{+12}$ ($Z=12, 13$ and $14$ respectively) is shown in figure 3.5. For the linear Stark effect to take place the quadratic Stark splitting must be much greater than the fine structure splitting in a given spectroscopic multiplet. The threshold for this in all three cases occurs at values above approximately $2 \times 10^8 \text{ V cm}^{-1}$ where
3.4 Plasma spectroscopy

Figure 3.4: The splitting of the Aluminium $^1S$, $^1P$ and $^1D$ states in the presence of an electric field.

Figure 3.5: Ratio of quadratic Stark splitting to the fine structure splitting for Magnesium, Aluminium and Silicon $^1P$ and $^1D$.

the ratio reaches values greater than 5 times the fine structure splitting. The intercombination transitions with triplet upper states $^3S$, $^3P$ and $^3D$ have more complex splittings and are therefore not included here.

When the external electric field is intense enough that the Stark splitting approaches and exceeds, for a given principal quantum number $n$, the fine structure splitting the linear model is applicable. The orbital quantum numbers can be separated into parabolic quantum numbers $n_1$ and $n_2$. The angular momentum quantum number $l$ is no longer used. The parabolic co-ordinates $n_1$ and $n_2$ are dependent upon the principal and magnetic quantum number where $n = n_1 + n_2 + ml + 1$. The electric field $F$ splits the upper $n$ and lower $n'$ states
into \((2n - 1)\) and \((2n' - 1)\) equidistant Stark sublevels separated in frequency units by,

\[
\omega_F(n) = \frac{3nq \hbar F}{2Zm_e e} \tag{3.9}
\]

Where \(q = n_1 - n_2\). The Stark sublevels of the upper and lower states have the designation \(\alpha \equiv (n, q, m_l)\), and \(\beta \equiv (n', q', m'_l)\) respectively. There is no selection rule with respect to the parabolic quantum numbers. However the transitions which involve a change of sign of \(n_1 - n_2\) are mostly weak. As usual the selection rule for the magnetic quantum numbers is \(\Delta m_l = 0\) for emission polarised parallel to the field whereas a change of \(\Delta m = \pm 1\) produces circular polarisation. There is no splitting of the state \(n = 1\) as \(q = n_1 - n_2 = 0\). Figures 3.6a and 3.6b show the linear Stark splitting upon a transition between \(n = 3 \rightarrow 1\) and the theoretical splitting diagram of the spectral lines. Transitions into excited lower states can give useful information about the density and temperature of plasmas. The splitting becomes more complex for transitions from higher values of \(n\) and to lower excited levels \(n' > 1\). For example the transition \(n = 3 \rightarrow 2\) contains 15 components as opposed to 5 for the simpler \(n = 3 \rightarrow 1\) as shown in figure 3.6a.

### 3.5 Field interactions

In this section the interaction between high frequency strong electric fields with atomic energy states is described. Two processes are outlined: large amplitude plasmons acting upon Stark split sublevels (Intra-Stark Spectroscopy), and the perturbation upon energy levels caused by a strong laser field. Both mechanisms produce characteristic spectroscopic effects which can be used to examine the characteristics of the plasma from which they are emitted.
3.5 Plasma spectroscopy

Figure 3.6: Linear Stark effect on the $n = 3 \rightarrow 1$ transition. The energy level scheme a) shows the splitting of the constituent states with respect to the parabolic and magnetic quantum numbers, and b) the theoretical splitting diagram of the spectral lines. The numbers against each line denote the relative frequency shift upon it.

3.5.1 Intra-Stark Spectroscopy

Intra-Stark Spectroscopy, as introduced in chapter 1, is the study of the dynamics and resonance between bound-bound transitions, low and high frequency electric fields. The resonance between each field produces particular structures within the emitted lines profile. Anticipating that the high frequency electric field is due to plasmons, or Langmuir waves, we adopt the expression “L-dip” for the structures. The theory of the L-dip effect was developed in a series of papers [30,31]. L-dip structures have been reported by various experimental groups using Z-pinches, gas-linear pinches, where electron densities inferred from L-dip positions in UV emission have agreed well with Thomson scattering results [30,35]. More recently X-ray transitions in long pulse ($t > 500$ ps) laser-produced plasma have been studied [36]. For spectral lines with higher photon energies the positions of the L-dips are closer to the central wavelength for a given electron density, therefore experimental observation of L-dips in X-ray line shapes requires instrumentation with high spectral and spatial resolution.
3.5 Plasma spectroscopy

Low frequency field sources include ion-acoustic waves, cyclotron motion or the Holtsmark field (as described in section 3.3). Large amplitude plasmons are the primary source of strong high frequency electric fields [34]. The radiating ions situated within an ion-microfield in a plasma are also subject to a perturbing field $E$. The total field is:

$$E(t) = F + E_0 \cos(\omega t)$$ (3.10)

where $F$ represents a low frequency quasistatic (with respect to the lifetime of the transition) field and $E_0 \cos(\omega t)$ is a monochromatic (single valued) high frequency field. The frequency of the field is the plasmon frequency $\omega$. An ensemble of radiators has a characteristic distribution $W(F)$ as a function of the absolute values of $F$. An ensemble of radiators will experience a quasistatic field with a Stark splitting of the upper and/or lower level of a transition. If this is equal to that of the frequency of the quasimonochromatic field, the system is said to be in a resonant state. Now the relation between the Stark splitting and the high frequency field can be written,

$$p\omega \approx \frac{3nhF}{2Zm_e}$$ (3.11)

where $p = 1, 2, 3, \ldots$ is the number of quanta (or plasmons) of the high frequency field involved in resonance. At intervals of $\Delta \omega\alpha = n^{-1}(nq - n'q')p\omega$ there is an abrupt intensity decrease at $\Delta \omega\alpha$ and an intensity increase either side of the resonance. This structure is called an L-dip. The location of the intensity increase relative to the unperturbed line centre $\lambda_0$ is dependent upon the amplitude of the quasimonochromatic electric field involved in the interaction. This is named the “half-width” and in frequency terms is,

$$\Delta \omega^{dip}_{1/2} \approx \left(\frac{3}{2}\right)^{1/2} \frac{n^2hE_0}{4m_e}$$ (3.12)

The depth of the dip is approximately equal to the distance between the two local maxima i.e. $2\Delta \omega^{dip}_{1/2}$. 
Figure 3.7: Calculation of L-dips due to a resonant interaction between a quasi-monochromatic and quasistatic field acting upon a spectral lineshape (solid red), and an unperturbed broadened lineshape (dotted blue) for a \( n = 3 \rightarrow 1 \) transition. Quadrupole interaction has been ignored.

Figure 3.7 shows an example of resonance upon a \( n = 3 \rightarrow 1 \) transition. Here the line shape is Lorentzian and \( \Delta \omega_{1/2} = 1.5 \times 10^{14} \text{ s}^{-1} \). Five dips each with two surrounding local maxima are found due to the interaction between the field and each of the \( 2n - 1 \) Stark sublevels. The polarisation of each dip structure is dependent on the magnetic quantum number of the upper and lower transition state, see figure 3.6b. In some cases where a dip is located at a portion of the lineshape with a large intensity gradient the peak (or maxima) nearest to the line centre is not as pronounced and instead a small increase in intensity results. In real spectra once the positions of the dips relative to the line centre are established and their half widths are measured then it is possible to infer the frequency, energy and amplitude of the plasmon involved in the interaction.

In high density plasmas, common to all laser solid interactions, the primary source of the low frequency quasistatic field are ion microfields. A quadrupole interaction between the radiator and the perturbing nearby ions of charge \( Z_p \) influences the energy Stark sublevels, therefore the position of the dips within the emitted spectral line will change accordingly. Taking into account that the oscillating high frequency field resonates with the upper Stark sublevels \( \alpha \) and
α + 1 the position of the dip relative to the line centre is [30],

\[ \Delta \omega_{\alpha+}(p) \approx n^{-1}(nq - n'q')p \omega + 2^{1/2}(27n^3ZZ_p \omega_a)^{-1/2}(p \omega)^{3/2}u_\alpha^+ \] (3.13)

\[ u_\alpha^+ \equiv n^2(n^2 - 6q^2 - 1) \]

\[ + 6n(2q + 1)(nq - n'q') - (n')^2[(n')^2 - 6(q')^2 - 1] \] (3.14)

\[ \omega_a \equiv me^4h^{-3} \approx 4.14 \times 10^{16}s^{-1} \] (3.15)

where \( \omega_a \) is the atomic frequency. The same is true for resonance between the oscillating field and the sublevels \( \alpha \) and \( \alpha - 1 \). In this case [30],

\[ \Delta \omega_{\alpha-}(p) \approx n^{-1}(nq - n'q')p \omega + 2^{1/2}(27n^3ZZ_p \omega_a)^{-1/2}(p \omega)^{3/2}u_\alpha^- \] (3.16)

\[ u_\alpha^- \equiv n^2(n^2 - 6q^2 - 1) \]

\[ + 6n(2q - 1)(nq - n'q') - (n')^2[(n')^2 - 6(q')^2 - 1] \] (3.17)

The L-dips are shifted by an amount dependent on the electron density. This shift is towards the blue/high energy portion of the lineshape. This amount in Amstrongs is,

\[ \delta \lambda(p) \approx 5.97 \times 10^{-15}p^{3/2}n_e^{3/4} \] (3.18)

Therefore at the point \( \Delta \omega_a = n^{-1}(nq - n'q')p \omega \) for a fixed number of interacting quanta two L-dips occur in place of a single L-dip at the positions \( \Delta \omega_{\alpha \pm} \) blue shifted by an amount equal to that found in equation 3.18. This high density phenomenon occurs in the same manner on the lower sublevel \( \beta \). Hence for the transition \( n = 3 \rightarrow 1 \) ten dips arise in the line profile whereas sixteen are found for the transition \( n = 3 \rightarrow 2 \). The number of dips will also increase with extra quanta \( p \) interacting with the radiator. In some cases dips due to the resonance between different sublevels will overlap. Therefore the dips observed in the spectrum may have extra structure; deeper dips with taller neighbouring maxima, or a series of suppressed features. As the complexity of the spectrum increases the confidence in determining the correct half widths of the dips and
3.6 Summary

In this chapter the radiative processes of a plasma was presented. Lineshape broadening by impact and ion motion was described. The splitting of energy components by the Stark effect was discussed. For helium-like ions the transition from a quasistatic Stark effect into a linear Stark effect is dependent on the element of the radiator and the transition of interest. For helium-like Al$^{11+}$ the linear Stark effect takes place for transitions with upper level principal quantum numbers $n = 3$ and 4 when the radiators experience ion-microfields greater than $2 \times 10^8$ V cm$^{-1}$. The resonant action between strong high frequency electric fields and the Stark split levels of the radiators is presented. The L-dip structures and their positions in lineshapes due to this interaction are described.
Chapter 4

The ASTRA experiment

4.1 Introduction

The purpose of the ASTRA experiment was to study the perturbations on atomic energy levels due to high frequency strong fields within a plasma. In order to achieve this a picosecond duration laser pulse is used to generate a plasma and to produce large amplitude plasma waves. In some cases a femtosecond duration laser pulse synchronised and delayed with respect to the picosecond pulse at adjustable time intervals after the picosecond pulse. The purpose of the femtosecond pulse was to illuminate the plasma and impose its characteristic laser induced electric field. In this chapter the layout of the ASTRA experiment and the X-ray diagnostics are presented. Time integrated, spatially resolved X-ray emission is presented. Modulations are observed in the lineshapes and are attributed to dynamic resonance between plasmons and the Stark split levels of the radiators.
4.2 Experiment design and diagnostics

4.2.1 The ASTRA laser system and target area

The experiment took place using the ASTRA laser system in Target Area 2 (TA2) at the Central Laser Facility, Rutherford Appleton Laboratory [64]. The Ti:Sapphire (Ti:S) laser provides pulses of 800 nm light with a minimum pulse duration of 40 fs at energies of up to 500 mJ at a rate of 2 Hz and maximum intensities on target of $10^{19}$ W cm$^{-2}$. Completed in 2000 its aim was to allow users to study femtosecond X-ray generation, ultra-fast probing of molecules, laser acceleration of particles, high harmonic generation and other high-intensity short pulse interactions. At the moment of writing this thesis ASTRA, due to its success, is being upgraded to a dual beam 0.5PW facility, a unique capability for a laser system (completion date Autumn 2007). Each beam will be independently focusable to $\sim 10^{22}$ W cm$^{-2}$ firing at a repetition rate of 2 shots per minute.

The seed pulses with a duration of 12 fs and an energy of 3 nJ are generated by a Ti:S oscillator system. Pulses are stretched to 530 ps successively amplified to 1.5 J. At this point a portion of the pulse is split compressed to 40 fs by a system of diffraction gratings. The average energy of the pulse into the target area is approximately 500 mJ with a shot rate of 2 Hz. The high repetition rate offered by ASTRA is an advantage in laser-plasma experiments as numerous shots can be taken in a day (>100).

The chamber in TA2 is cuboid in shape with a length, width and height of 1.04 m, 0.9 m and 0.6 m respectively and wall thickness of 0.075 m. The primary diagnostic within the chamber, a highly dispersive toroidal spectrometer (HDTES), was positioned to disperse time-integrated spatially resolved aluminium He-like K-shell spectra onto a large area CCD camera. Due to the nature and geometry of the diagnostic an extension was needed to contain the CCD camera.
4.2 The ASTRA experiment

4.2.2 Experimental layout

A layout of the experimental setup is presented in figure 4.1a and 4.1b showing side and top views respectively. The diagnostic layout including crystal spectrometer and CCD is shown. This layout was designed to accommodate two experiments. Equipment for the second experiment is shown in 4.1. The picosecond laser pulse was stretched to a full width half maximum (FWHM) duration of 3.4 ps whilst the femtosecond pulse duration was 35 fs. Both pulses contained an energy of 200 mJ and were operated with a repetition of 2 Hz. In some instances
4.2 The ASTRA experiment

Figure 4.2: Illustration of the laser-target geometry. The target is illuminated after the OAP focus.

<table>
<thead>
<tr>
<th>Spot size (µm)</th>
<th>Picosecond pulse (W cm(^{-2}))</th>
<th>Femtosecond pulse (W cm(^{-2}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>15</td>
<td>3\times10^{16}</td>
<td>3\times10^{18}</td>
</tr>
<tr>
<td>50</td>
<td>3\times10^{15}</td>
<td>3\times10^{17}</td>
</tr>
<tr>
<td>100</td>
<td>8\times10^{14}</td>
<td>7\times10^{16}</td>
</tr>
</tbody>
</table>

Table 4.1: The on target intensities of the primary and secondary pulses as a function of spot size.

only the primary or secondary pulse was used. The laser had a relatively low pulse contrast, with a pedestal to peak intensity ratio of 1:10\(^3\) at 100 ps before the peak of the pulse. The contrast was not measured. Previous CLF analysis of the laser pulse is referred to [65]. The time delay of the femtosecond pulse relative to the picosecond was varied between 50 ps and 400 ps to illuminate the plasma created by the picosecond pulse at different stages of its evolution. The laser pulse travels along the highlighted path as shown in figure 4.1 and is focused by a f/6.5 off-axis parabola (OAP) onto the target with a 10° angle of incidence. By defocussing the OAP the focal spot radius was varied between 15 µm and 100 µm after best focus as presented in figure 4.2. This produced a range of intensities on target as shown in table 4.1 whilst keeping the energy in the beam constant.

Targets consisted of 500 µm diameter 5 cm long aluminium wires. These were used to take advantage of ASTRA’s high repetition rate, allowing multiple shots
4.2 The ASTRA experiment

to be taken in a short period of time. Maximum emission occurs at the target surface where the density of emitters is greatest. Here the radiators are perturbed primarily by Stark broadening mechanisms. Emission reduces further away from the target surface due to the decrease in plasma density. To accommodate for the anticipated low emission in the laser plume it was necessary to integrate a single data set over a number of consecutive shots. This was achieved by translating the wire along the focal plane of the laser. Each shot was made on a smooth portion of the target with an intershot spacing of approximately 0.5 mm. Operating at 2 Hz it was possible to accumulate 100 shots in one data set taking approximately 50 seconds to complete.

The target mount consisted of an aluminium holder connected to a motorised XYZ stage. The mount needed to translate the target through its entire length in the laser focal spot. A glass rod with equal dimensions as the aluminium target was used for target alignment. By obscuring the laser focal spot with the glass rod the XYZ stage was fixed into the correct position and angle. The target was fixed to the mount securely to ensure the target did not buckle due to laser illumination. Debris shields were fitted to the target mount and to the diagnostics to halt damage.

4.2.3 X-ray diagnostics

Time integrated, spatially resolved X-ray emission was recorded using a highly dispersive toroidal spectrometer HDTS (100) fitted with a quartz crystal with a 2d lattice spacing of 8.5009 Å. The crystal horizontal and vertical bending radii are 1404.55 mm and 312.80 mm respectively. Crystal manufacturing precision is greater than 1 µm. The emission spectrum was resolved parallel to the target surface with spatial resolution perpendicular to the target surface. This line of sight allows for spatial resolution of the plasma parallel to the expansion axis. Each cross-section will contain emission from approximately uniform
plasma conditions. The spectral emission can be analysed and compared directly to simulations as a function of height above the target surface. This is unusual and possibly unique. The crystal was protected from target debris with two layers of aluminised Mylar foil consisting of 14 nm of aluminium on one face of a 7 μm thick film of Mylar (C_{10}H_{8}O_{4}). An HDTs was used due to its high dispersion, luminosity and source size independent spatial resolution [66]. It provides a spectral resolution greater than $\lambda/\Delta\lambda = 8000$. The large resolving power is needed to distinguish any small spectral features in the lineshape. The spatial resolution was limited by the HDTs magnification and CCD pixel size to 3μm.

Figure 4.3: An illustration of the target and diagnostic setup.

The energy dispersed X-ray emission from the plasma was recorded onto a large area scientific-grade deep depletion, front illuminated e2V CCD42-90 camera and 16-bit amplifier sensitive to 6 electron-hole pairs manufactured by XCAM [67]. X-ray photons interact with a silicon layer called the “depletion layer”. Photons are absorbed by valency electrons in this layer which are promoted across the band gap to the conduction band. The number of promoted electrons is dependent on the energy of the absorbed photon. The electrons are collected by an array of capacitors which are read out and converted into a two
4.2 The ASTRA experiment

<table>
<thead>
<tr>
<th>Spectral line</th>
<th>Configuration</th>
<th>Term</th>
<th>Energy (eV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>He β</td>
<td>1s3p → 1s2</td>
<td>1P1</td>
<td>1S0</td>
</tr>
<tr>
<td>He β IC</td>
<td>1s3p → 1s2</td>
<td>3P1</td>
<td>1S0</td>
</tr>
<tr>
<td>Li-like satellites</td>
<td>1s3lnl' → 1s2nl</td>
<td>multiple</td>
<td>1849.37...1860.75</td>
</tr>
<tr>
<td>He γ</td>
<td>1s4p → 1s2</td>
<td>1P1</td>
<td>1S0</td>
</tr>
<tr>
<td>He δ</td>
<td>1s5p → 1s2</td>
<td>1P1</td>
<td>1S0</td>
</tr>
<tr>
<td>He ε</td>
<td>1s6p → 1s2</td>
<td>1P1</td>
<td>1S0</td>
</tr>
</tbody>
</table>

Table 4.2: Al$^{+11}$ K-shell spectral lines and Li-like satellite lines recorded by the HDTS.

dimensional image. In a front-illuminated device the capacitor array faces the incoming radiation. Therefore photons must travel through the capacitor array before arriving at the photosensitive deep depletion layer. The deep depletion layer ensures photon absorption in a single pixel and limits charge migration to neighbouring pixels thereby reducing “split-events”.

The CCD chip area is 62 mm × 28 mm containing 4608 × 2048 pixels with a size of 13.5 × 13.5 µm. Prior to recording data the CCD chip is cooled to -60 °C by a multistage Peltier cooler. Quantum efficiency at Al$^{+11}$ K-shell energies is approximately 50%, with a dynamic range of 300 where each absorbed He β (1.9 keV) photon generating about 530 electron-hole pairs with an energy resolution of 25 eV. To ensure the CCD was light tight and protected from hard X-rays two additional layers of aluminised Mylar were used to cover the CCD. Furthermore to protect the CCD from hot electron impacts it was separated from the main target chamber with slits, lead screens and a 1 T bending magnet as shown in figure 4.3. A CCD was chosen as it is by far the most useful instrument to record weak X-ray signals compared to other detectors such as X-ray film and imaging plates [4]. CCDs have superior signal-to-noise ratios and are more practical in high repetition rate experiments where the shot-to-shot time is not limited by image development and detector replacement.

Due to the high dispersion properties of the HDTS it was necessary to use two positions of the crystal and detector; one to record Al$^{+11}$ K-shell He β,
intercombination line (IC) and neighbouring Li-like satellites (energy range 1890 - 1830 eV), and another to record He $\gamma$, $\delta$ and $\epsilon$ (energy range 2030 - 1935 eV). Details are shown in table 4.2. Great care was needed to position the crystal and CCD correctly. The HDTS and the CCD were separately mounted onto movable stages to allow for precise alignment. The positioning was better than 0.1 mm. Details are shown in table 4.3. The target-to-crystal and crystal-to-CCD distances determine the spectral range, the magnification factor of the plasma and hence the spatial resolution. Therefore the magnifications of the two setups are 4.5 and 3.8 and the spatial resolutions limited by the 13.5 $\mu$m pixel size are 3.0 $\mu$m and 3.6 $\mu$m.

### 4.3 Experimental data

Data collected by the CCD is processed firstly by identifying the recorded spectral lines and secondly accounting for the imaging properties of the spectrometer. This includes converting the counts per pixel into the number of detected photons per micron squared using the electron-hole pair conversion and the quantum efficiency of the CCD for the Al $^{+11}$ K-shell energy range [68]. Figure 4.4a shows a typical He $\beta$ recording with Li-like satellites identified and figure 4.4b presents a typical He $\gamma$, $\delta$ and $\epsilon$ recording. The dispersed photon energy is shown along the horizontal axis, whereas spatially resolved emission is shown along the vertical axis. The vertical axis is expanded. Spectral information is detected to 200 $\mu$m above the target surface. Satellites are seen on the low energy wing for all spectral lines. This accounts for the asymmetry of the line shape near the target.

\[
\begin{array}{|c|c|c|}
\hline
\text{Spectral line} & \text{target-to-crystal} & \text{crystal-to-CCD} \\
& (\text{mm}) & (\text{mm}) \\
\hline
\text{He $\beta$ $\rightarrow$ Li-like satellites} & 245.6 & 1095.1 \\
\text{He $\gamma$ $\rightarrow$ He $\epsilon$} & 271.7 & 1023.3 \\
\hline
\end{array}
\]

Table 4.3: Positions of the HDTs and the CCD.
4.3 The ASTRA experiment

surface. Close to the initial target surface the He-like resonance lines and Li-like dielectronic satellites are clearly resolved in figure 4.4a. These satellites are strongly correlated to excited $1s n l 3 l'$ states populated by collisional processes at high plasma densities, where $l$ and $l'$ are the orbital angular momentum states of the excited levels [69].

Transverse cross sections of the data are taken to examine the evolution of line shapes as a function of distance above the target surface. Cross sections of the data from the two spectrometer positions are displayed in figure 4.5a and 4.5b with a spatial resolution of 6 µm and 7.2 µm (averaged over 2 pixels to increase signal-to-noise ratio $S/N$, see section 4.3.2) respectively for heights above target surface between 0 µm and 100 µm. It is evident that close to the target surface the linewidths are at their broadest. Further from the target surface the lines narrow. This reduction in the line width can be attributed to Stark broadening with a small (approximately 10 %) contribution from Doppler broadening and the decrease in electron densities further from the target surface.

4.3.1 Data smoothing

One feature that is evident within the cross sections, particularly emission from near target surface, is the abundance of high frequency noise on the lineshape. This noise is not due to spectral resolution limits or the CCD structure. Along the spectral direction the CCD oversamples the spectrum i.e. the rocking curve of the crystal is a few of pixels wide. The source of the noise is probably due to a combination of the readout noise of the CCD and the high background counts with respect to the plasma emission. To reduce this noise, improve the signal-to-noise ratio and gain access to the underlying lineshape structure it is necessary to employ smoothing algorithms. A fast-Fourier transform (FFT) smoothing method is used [70]. This is accomplished by removing Fourier components above a defined threshold frequency $1/n_{FFT} \Delta t$ where $n_{FFT}$ is the number of data points.
4.3 The ASTRA experiment

Figure 4.4: An example of raw data for the two spectrometer positions; a) He $\beta$ and Li-like satellites, b) He $\gamma$, $\delta$ and $\epsilon$.

Figure 4.5: Spatially resolved cross sections of a) He $\beta$ and Li-like satellites, and b) He $\gamma$, $\delta$ and $\epsilon$. 
considered at one time and $\Delta t$ is the spacing between two adjacent data points. Frequency space in this case has units of eV$^{-1}$. Larger values of $n_{FFT}$ result in lower cutoff frequencies, and thus a higher degree of data smoothing. It is superior to other forms of smoothing, such as adjacent averaging, where averaging generally “washes out” the peak height and width of the features in data. It is imperative that $n_{FFT}$ is not sufficiently high (low frequency cutoff) as to smooth out finer features associated with atomic perturbations altogether. Figure 4.6 shows the effect on data when $n_{FFT}$ increased from low to high values for the He $\beta$ spectral line spatially resolved at target surface. Even at low values of $n_{FFT}$ ($n_{FFT} = 2$ in figure 4.6) where the frequency cutoff is quite high ($\approx 30$ eV$^{-1}$) the high frequency noise is reduced and the underlying structure is revealed on the line profile. As $n_{FFT}$ is increased noise is suppressed and features become more evident. At higher values of $n_{FFT}$ ($> 12$; low frequency cut off $< 5$ eV$^{-1}$) the features are smoothed out completely making analysis difficult. Furthermore the structure on the lineshape is not an artifact of the FFT smoothing process. The structure is found using other smoothing mechanisms such as Binomial smoothing and Savitzky-Golay smoothing. In all subsequent cross sections $n_{FFT} = 8$ as it is a good compromise between noise suppression and modulation identification.

The modulations are interesting. They are observed in the region close to
the target surface (within approximately 30 \( \mu m \)). At greater heights above the target surface the modulations reduce in number and magnitude. They occur in both picosecond only and double pulse data. Modulations are strongest and more numerous in picosecond only data. In some double pulse cases the modulations are not observed. This is not due to greater signal-to-noise ratio as a similar number of photons are collected in each data set. The line shape modulations are discussed in section 4.3.3.

4.3.2 Signal-to-noise ratio

To assert a level of confidence on the origin of the features the signal-to-noise ratio must be considered. If the amplitude of the features is below the noise level then it is more than likely that the features are random fluctuations in the line profile. Otherwise the features must be considered to be a genuine artifact of atomic energy level perturbations of the radiators within the plasma. The signal-to-noise ratio \( S/N \) of a given detector area is defined by equation 4.1 as

\[
\frac{S}{N} = \frac{N_s}{\sqrt{N_s + n_{pix}(N_S + N_D + N_R^2)}}
\]  

(4.1)

where \( N_s \) is the total number of photons, \( N_S \) is the average number of background counts per pixel, \( N_D \) is the dark-signal per pixel (counts created due to thermal energy within the CCD silicon lattice), \( N_R \) is the read out noise per pixel (generated by the CCD amplifier), and \( n_{pix} \) is the number of pixels within the detector area under consideration [71]. By increasing the sampling area over which emission is averaged a larger number photons will be counted. Therefore if the sampling area is increased the signal-to-noise ratio also increases as \( S/N \propto N_s/\sqrt{N_s} = N_s^{1/2} \). A cross section with a 2 pixel height and covering an energy range of 1860 - 1875 eV (He \( \beta \) ) and dispersion of 0.01709 eV pixel\(^{-1}\) contains 1756 pixels. For the unsmoothed line shape used for figure 4.6 there
4.3 The ASTRA experiment

is a total number of photons $N_s = 10832$, average background count per pixel $N_S \approx 8$, a read out noise of $N_R = 3$, and a negligible dark signal. Therefore, accounting for the 2 pixel averaging, the signal-to-noise ratio $S/N \approx 300$. In other words noise on the line profile is a factor of $1/300^{th}$ of the modulations. Measuring from peak to trough the features are approximately 80 times greater in magnitude in comparison with noise level. Therefore the features are a real artifact of the atomic transition process. This is strengthened by the fact that some data recorded using the same experimental setup but with different laser parameters does not contain this particular lineshape structure.

4.3.3 Spatial sampling

Due to the extremely high spatial and spectral resolution of the HDTS-CCD combination it is possible to record and identify very small perturbations on lineshapes. The spatial resolution of the spectrometer is limited by the pixel size of the CCD. By averaging emission over a number of pixels the signal-to-noise ratio increases. However in doing so any features within the lineshape which are spatially dependent will be obscured as the spatial resolution decreases. The laser pedestal itself is intense enough to create a plume of plasma extending out tens of microns above the surface through which the picosecond pulse travels. The pedestal does not ionise the aluminium to K-shell. By averaging over larger numbers of pixels information is lost such as the electron density and electric field strength for a particular given location. Due to the time-integration of the data, averaging will take place over regions where the plasma created by the low contrast pedestal had not extended when the picosecond pulse arrived. Hence cross sections with a large sampling area will contain both emission from locations where laser-plasma coupling has taken place and expanded plasma due to the primary pulse. Therefore the number of pixels used to average the emission should be kept to a minimum. Figure 4.7 shows an example of the He $\beta$ lines.
profile taken at the target surface for a different number of pixels over which the emission is averaged. It is still possible to identify peak and dip features throughout the line profile for large numbers of pixel averaging in the vertical spatial direction in figure 4.4. It is clear that the number and depth of features reduces with decreasing spatial resolution. This affects the analysis of the data as the position and structure of the modulations are sensitive to plasma parameters thus the associated analysis and errors will be dependent on the sampling area. In all subsequent data analysis cross sections are made over a two pixel sampling area.

The modulations are not an artifact of the CCD structure. The modulations are observed for distances up to 30 \( \mu m \) above the surface. Furthermore the modulations are not observed for some double pulse shots. Neighbouring satellites to the He \( \beta \) line are found at energies less than 1865 eV and do not contribute to the structure of the line shape. Therefore the modulations are probably due to the interaction between plasmons (Langmuir-waves) supported by the plasma and the radiators. The modulations will be referred to as L-dips (Langmuir-dips) and the position and structure of the modulations will be analysed using Intra-Stark theory [30].
4.4 Data comparison

In this section qualitative and quantitative comparisons are made between picosecond only pulse data and double pulse data for a range of intensities and time delays. Initially a description and analysis of the single primary pulse He \( \beta \) data for different focal spot intensities is made (due to the relatively reduced complexity of it’s Stark sublevel structure as discussed in section 3.5.1). Electron densities and high frequency electric field spatially resolved measurements are extracted from the L-dip structure. Dual primary/secondary pulse He \( \beta \) data is then compared to single pulse He \( \beta \) data. Finally the He \( \gamma, \delta \) and \( \epsilon \) transitions will be analysed in the same manner.

4.4.1 Picosecond pulse only He \( \beta \) data

Figure 4.8 shows cross sections from data for an intensity of \( 3 \times 10^{16} \) W cm\(^{-2} \) for different spatial locations above the initial target surface. It is evident that L-dip features are more numerous and deeper in the cross sections close to target surface. Further from target surface the L-dips narrow, reduce in number and become less deep until approximately 24 \( \mu \)m where features are not observed in the line profile. This is possibly due to the spatial dependence of the electric field strength. Furthermore, as seen in figure 4.8 indicated by dotted vertical lines, a number of L-dips occur at the same (or very close) spectral position for different spatial locations. This suggests that a group of radiators within the plasma is subject to a narrow band single high frequency field for all spatial locations up to approximately 20 - 30 \( \mu \)m above the target surface. The bandwidth of the high frequency electric field as seen by the radiators from target surface to this position depends on the density and the density gradient of the plasma \( L_n \). L-dip positions are sensitive to spatial location. The maximum electric field will be reduced at the extremes of the plasma expansion due to lower densities. Collisional damping is
also increased at the extremes of plasma expansion due to greater temperatures. L-dip width, depth and number will also be sensitive to spatial location. The hydrodynamic and laser-plasma coupling mechanism simulations are covered in section 2.

Figures 4.9a and 4.9b show cross sections from data between the target surface and 24 µm above target surface for intensities of $3 \times 10^{15}$ W cm$^{-2}$ and $7 \times 10^{14}$ W cm$^{-2}$ respectively (laser focal spot sizes of 50 µm and 100 µm). It is evident that L-dip formations are found in these cross sections at heights up to, and greater than, 24 µm above target surface as in the case of data extracted using an intensity of $3 \times 10^{16}$ W cm$^{-2}$. Furthermore the number of L-dips and relative depth decreases with height, i.e. the number of quanta $p$ that interact with and perturb the Stark sublevels decreases. This suggests the strength of the high frequency electric field associated with plasma waves diminishes as a function of height for all single pulse data.

Extracting electron density and electric field measurements from cross sections of the data is not a simple process. Firstly the complexity of the quanta/sublevel interaction in the case of He β produces two L-dips $\Delta\omega_{\alpha\pm}$ per sublevel due to the quadrapole interaction between ions in high density plasma as discussed in
Figure 4.9: A comparison of He $\beta$ cross sections for different spatial locations between 0 $\mu$m and 24 $\mu$m for a laser focal spot size of a) 50 $\mu$m and b) 100 $\mu$m. Heights above the target surface are 0$\mu$m (solid line), 12$\mu$m (dashed line), 18$\mu$m (dotted line) and 24$\mu$m (dash-dot line) are presented.
chapter 3 section 3.5.1. The L-dips are blue shifted an amount dependent on the number of quanta \( p \) interacting with the sublevels. Also the L-dips resulting from \( p \) quanta interacting with one set of radiators may constructively or destructively interfere with L-dips resulting from for example \( p + 1 \) with another set of radiators. Two L-dips which are sufficiently close together may create a “pseudo” dip between the two hills adjacent to them. The plasma parameters are extracted from data by calculating L-dip positions using a number of electron densities and number of quanta involved in the interaction.

For the He \( \beta \) cross section shown in figure 4.10 taken at target surface for an intensity of \( 3.3 \times 10^{16} \) W cm\(^{-2} \) the frequency of the plasmons involved in the interaction is found to be \( 9.0(\pm 1.0) \times 10^{14} \) rad s\(^{-1} \) which corresponds to an electron density of \( n_e = 2.0(\pm 0.4) \times 10^{20} \) cm\(^{-3} \). The vertical dashed lines indicate \( \Delta \omega_{\alpha+} \) L-dip positions for clarity as \( \Delta \omega_{\alpha-} \) L-dips in some cases overlap, particularly for low \( p \). The electric field of the plasmons involved in the interaction inferred from data is found to be \( E_0 = 1 \times 10^8 \) V cm\(^{-1} \) and the number of quanta is found to be \( p = 4 \). Table 4.4 shows the positions of \( \Delta \omega_{\alpha\pm} \) L-dips in energy units (eV), the associated value of \( q = n_1 - n_2 \) from chapter 3 section 3.5.1, quanta involved in the interaction (excluding \( p = 4 \)) and the polarisation of the dip. It is clear that increasing the number of quanta involved in the interaction leads to a more complex spectrum of L-dips.

At 12 \( \mu \)m above the surface the density inferred from the L-dip positions decreases slightly to \( 1.8(\pm 0.4) \times 10^{20} \) cm\(^{-3} \). The high frequency electric field remains at \( E_0 = 1(\pm 0.2) \times 10^8 \) V cm\(^{-1} \). However it must be stated that this value is inferred from an average of the L-dip half widths and therefore the percentage error is significantly higher than that associated with an electron density measurement inferred from L-dip positions. Above 12 \( \mu \)m L-dip features are not observed. This suggests either that the electric field of the plasmons has diminished to a point where it is much less than the ion microfield \( E_0 << F_0 \) and the perturbation does
4.4 The ASTRA experiment

Figure 4.10: The He $\beta$ line cross section at target surface for a laser spot size of 15 $\mu$m. The vertical lines indicate L-dips theoretical L-dips for $n_e = 1.8 \times 10^{20}$ cm$^{-3}$. The positions of the dips are presented in table 4.4.

<table>
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<th>$\Delta \omega_{\alpha+}$ (eV)</th>
<th>$\Delta \omega_{\alpha-}$ (eV)</th>
<th>q</th>
<th>polarisation</th>
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<td>1873.251</td>
<td>2</td>
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</table>

Table 4.4: Calculated L-dip positions using the model of chapter 3.5.1 on the He $\beta$ lineshape cross section taken at target surface for an intensity of $3 \times 10^{16}$ W cm$^{-2}$. 
not take place, or that the resolution of the instruments is not high enough to distinguish the features at lower electron densities.

A similar trend is also found in data for lower intensities of $3 \times 10^{15}$ W cm$^{-2}$ (50 µm spot size) and $8 \times 10^{14}$ W cm$^{-2}$ (100 µm spot size). For an intensity of $3 \times 10^{15}$ W cm$^{-2}$ the density at target surface inferred from L-dip positions is $1.8(\pm 0.4) \times 10^{20}$ cm$^{-3}$ decreasing to $1.3(\pm 0.4) \times 10^{20}$ cm$^{-3}$ at higher spatial locations. In this case L-dips are observable up to 24 µm. The electric field also decreases from $E_0 = 1(\pm 0.2) \times 10^8$ V cm$^{-1}$ at target surface to $9(\pm 0.2) \times 10^7$ V cm$^{-1}$ at 24 µm. Furthermore the number of quanta involved in the interaction decreases from $p = 4$ to $p = 3$. For an intensity of $8 \times 10^{14}$ W cm$^{-2}$ the density reduces from $2.6(\pm 0.4) \times 10^{20}$ cm$^{-3}$ at target surface to $1.7(\pm 0.4) \times 10^{20}$ cm$^{-3}$ at 18 µm above the target surface. Above this location L-dips are not observed. The electric field reduces from $E_0 = 1(\pm 0.1) \times 10^8$ V cm$^{-1}$ at target surface to $9(\pm 0.2) \times 10^7$ V cm$^{-1}$ at 18 µm. The number of quanta $p$ decreases from 4 at target surface to 2 at 18 µm above target surface. An explanation for why L-dip features are seen at greater distances above the target surface for lower intensities may be due to an inaccuracy in cross section extraction from raw data. It is likely the distance above target surface has an error of 1 to 2 pixels, or 3 to 6 µm with respect to the plasma.

### 4.4.2 Double pulse He $\beta$ data

In addition data was recorded using a double pulse for comparison purposes consisting of a picosecond 3.4 ps pulse and a femtosecond 35 fs pulse separated by a controllable time delay. Data using time delays of 50 ps to 400 ps were recorded; however due to experimental time constraints full data sets could not be recorded in some instances. For a focal spot size of 15 µm the time delays used were 50, 100 and 200 ps; for 50 µm the time delays were 50, 200 and 400 ps, and for 100 µm the time delays were 100 and 200 ps. However data is comparable
Figure 4.11: A comparison of He $\beta$ cross sections for different spatial locations between 0 $\mu$m and 24 $\mu$m for a laser focal spot size of 100 $\mu$m and a time delay of 100 ps. Heights above the target surface are 0 $\mu$m (solid line), 12$\mu$m (dashed line), 18$\mu$m (dotted line) and 24$\mu$m (dash-dot line) are presented.

between all focal spot data sets for at least one time delay. The main feature of the double pulse data compared to the single pulse data is the greatly reduced number of L-dips on the spectral line shape. Most of the L-dips are found from cross sections at the target surface. In some cases L-dips do not appear at all, this is particularly the case for shorter time delays (50 ps and 100 ps). As an example figure 4.11 shows cross sections taken from data using a 100 $\mu$m focal spot for a 100 ps time delay between the primary and secondary pulse. It is quite clear there is a lack of L-dip structures as compared with cross sections from 15 $\mu$m spot size data. However a trend that most of the data shows is that L-dips reappear, have a more pronounced structure and a greater number of quanta $p$ for longer time delays. There may be a number of explanations for this trend. Firstly the lack of L-dip structures for all double pulse data for a spot size of 100 $\mu$m may be due to the lateral spatial extent of the plasma. Cross sections taken through greater cross sectional areas will contain non-uniform plasma parameters. Therefore the high frequency field associated with plasmons contained within the plasma may be sufficiently broadband as to “smear out” any L-dip structures. However this does not seem to be the case for single pulse data for the same focal spot size. Therefore to explain the general trend in the data we must turn to the
### Table 4.5: Experimental electron densities and high frequency electric field strengths inferred from He $\beta$ data. Note all double pulse 100 $\mu$m spot size data do not contain L-dips.

<table>
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<th>Height ($\mu$m)</th>
<th>$n_e$ (cm$^{-3}$)</th>
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<th>$p$</th>
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The hydrodynamics of the secondary pulse interacting with the plasma formed by the primary pulse.

Table 4.5 presents an overview of all of the data including single pulse regarding electron density $n_e$, electric field $E_0$ and number of quanta $p$ inferred from data as a function of spatial location. The majority of the inferred electric field strengths is found to be $\approx 1 \times 10^8$ V cm$^{-1}$. This suggests that during the laser-plasma interaction plasma waves are amplified to a maximum value as dictated by wave breaking mechanisms and particle trapping, both of which are dependent upon the local electron density. The inferred electron density for each line is approximately equal for all lines with a small deviation from the mean. Therefore the maximum electric field will be limited to approximately the same level.
4.4 The ASTRA experiment

Figure 4.12: A cross section taken at target surface using a 50 µm spot size and single pulse only.

4.4.3 He $\gamma$, $\delta$ and $\epsilon$ data

Figure 4.12 shows He $\gamma$, $\delta$ and $\epsilon$ line shapes from the same cross section taken at target surface from 50 µm spot size primary pulse only data and figures 4.13a, b and c show each line shape in greater detail for each transition respectively. Unfortunately the He $\delta$ and $\epsilon$ lines noise level is a significant fraction of the signal. This is not the case for the He $\beta$ and $\gamma$ lines. Furthermore the complicated spectrum of He $\delta$ and $\epsilon$ adds further conclusive analysis complications with 11 and 13 sub-levels respectively each creating two L-dips in the line shapes for a single quantum interaction. Therefore only the He $\gamma$ lines will be analysed where 9 sub-levels create 17 L-dips in the line shape for a single quantum interaction. For low densities as inferred from the He $\beta$ analysis the two L-dips from one sublevel $\Delta \omega_{\alpha\pm}$ are very close and re-enforce one another making the L-dips much easier to discern.

Table 4.6 presents electron densities, electric field strengths and number of quanta involved in the interaction inferred from He $\gamma$ data as a function of height above the target surface. A similar trend is found compared with the He $\beta$ single and double pulse analysis. L-dips are found primarily for single pulses in cross sections taken at the target surface and for double pulse data fewer cross sections are found containing L-dips. There are some interesting differences between the
Figure 4.13: A cross section taken at target surface using a 50 μm spot size and single pulse only.
4.5 The ASTRA experiment

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Table 4.6: Electron densities and high frequency electric field strengths inferred from He \( \gamma \) data. Note all double pulse 100 µm spot size data do not contain L-dips.

He \( \beta \) and \( \gamma \) lines; the inferred electric field strength is on average lower and one quanta is involved in the interaction in the case of the He \( \gamma \) line. This observation has a linked explanation. The maximum rate of the He \( \gamma \) transition from \( n = 4 \rightarrow 1 \) may occur at a time when the plasmons have reduced in amplitude due to Landau damping. For the He \( \gamma \) L-dips the lower electric field strength inferred from data reduces the probability of a multiquantum interaction. This shall be covered in detail in the following chapters.

### 4.5 Summary

In this chapter the experimental details of the ASTRA experiment were presented and discussed. The X-ray spectrometer and CCD detection system were shown to be suitable diagnostics for this study offering excellent spatial and spectral resolving powers. Signal smoothing and signal-to-noise ratio were discussed to observe and verify the presence of modulations within the He \( \beta \) and He \( \gamma \) line profiles. The Electron densities, electric field strengths and number of interaction...
quanta were inferred from the L-dip structures. Discussion of the results was given with reference to the hydrodynamic and atomic kinetic parameters of the plasma. In the following chapters this discussion will be extended using the aid of laser-plasma simulations.
Chapter 5

Hydrodynamics simulations

5.1 Introduction

In this chapter hydrodynamic simulations are performed to support the interpretation of the experimental data. The simulations focus on the region close to the target surface. This choice is motivated by the experimental data and the observations of modulations, particularly in the He $\beta$ line shapes. An understanding of the He $\beta$ emission is gained by studying the excited $n = 3$ populations. In this section hydrodynamic simulations of the laser plasma interaction using the code EHYBRID [72–74] are discussed. This code allows calculations of electron and ion temperatures due to lateral heat conduction. This characteristic has been studied and compared to other hydrodynamics codes in planar expansion regimes [66].

Firstly a brief description of the code is given, then the simulation results are presented and discussed. Simulations using a high contrast picosecond pulse are compared to simulations using low contrast picosecond pulse. This is to examine the role of a pre-formed plasma in the generation of plasmons. The spatial and temporal evolution of hydrodynamic parameters such as the electron density,
temperature and the average. Ionisation is analysed, and then discussed in terms of the creation and lifetimes of parametric instabilities. The population density ions in the $n = 3$ and 4 levels, corresponding to the upper levels of the He $\beta$ and $\gamma$ transitions respectively, are studied.

Double pulse simulations are not performed as the femtosecond pulse will generate high electron oscillation velocities. This cannot be modelled in hydrodynamical simulations. In this case a Particle-in-cell (PIC) code must be used. This is beyond the scope of this thesis.

The presence, growth and damping of plasmons amplified by the two plasmon decay and stimulated Raman scattering parametric instabilities are calculated by post-processing the output data from EHYBRID. The electric field strengths associated with the amplified plasmons are compared to those inferred from data. The temporal evolution of the population densities and plasmon electric field strengths are compared. This is to study the interaction between the upper states of the He $\beta$ and $\gamma$ transitions and the high frequency electric fields associated with the plasmons. This is needed to compare experimental data to a complete simulation.

5.2 Ehybrid

EHYBRID is a 1.5 dimension Lagrangian hydrodynamic code designed for modelling the creation and expansion of a laser produced plasma. The half dimension is used to allow the plasma to expand laterally from the target. This 2D expansion influences plasma temperatures due to an increased rate of cooling. Targets in the code consist of one or more layers of material with a thickness and width chosen by the user. Each layer is divided into a number of cells of thickness specified by the user. For single element targets an automatic quadratic zoning of cells can be implemented. During a simulation, cells reaching temperatures higher than a
value as specified by the user are located. The cells are then re-zoned quadratically so that the hotter cells occupy $5/6$ of the original target thickness and the simulation is restarted. This is to maximise resolution in regions where strong laser absorption occurs.

Multiple laser pulses can be used in the simulation. Pulses can have a different wavelength, pulse duration, energy and focal spot size. Laser energy is absorbed into the plasma by an energy dump at critical density dependent on plasma reflectivity and via inverse bremsstrahlung, mechanisms discussed in previous sections. The pulse shape can be specified by the user to be Gaussian or trapezoidal in shape. The pulse illuminates the target with a Gaussian spatial distribution of intensity dependent on the focal spot size. The focal spot size can be specified to be different from the target width.

A flux limited Spitzer conductivity model is used to control thermal conduction in the plasma. The flux limit, normally set to 0.1, is set by the user. The plasma is treated using a two fluid model where the Maxwellian electron and ion velocity distributions are considered separately and characterised with different temperatures. The ionisation of the plasma is treated with a full time-dependent ionisation balance with a collisional-radiative treatment with an arbitrary number of atomic levels for specified ionisation stages specified by the user or generated by the code. The equation of state of the plasma incorporates contributions from the nuclear and the free electrons. The nuclear component is user specified to be either an ideal gas with an arbitrary adiabatic constant or a full equation of state using the CHART D formulation as developed by Thomson and Lauson [75]. An ideal gas or Thomas-Fermi model can be chosen to model the free electrons. An ideal gas can approximate low density plasmas, however as densities increase electron degeneracy must be considered. The Thomas-Fermi model describes high density situations with greater accuracy.
5.3 Simulation parameters

To simulate the experiment and understand the role of pedestal to peak contrast, two pulses are modelled. The first is a high contrast pulse, this is a Gaussian shape 3.4 ps FWHM pulse. The second is a low contrast pulse, this has a 3.4 ps FWHM pulse with a pedestal set at $10^{-3}$ of the peak intensity. This second pulse more closely represents the experimental beam. The contrast was determined by calculations [65]. In all subsequent simulations time is measured from the peak intensity of the picosecond pulse. Simulations are carried out for the three laser intensities corresponding to the focal spot sizes of 15 $\mu$m, 50 $\mu$m and 100 $\mu$m.

The simulated target consists of a single layer of aluminium, 200 $\mu$m thick divided into 98 cells. The target width is altered for each focal spot size so that more than 99% of the Gaussian beam irradiates the target. The single layer is used in conjunction with the quadratic re-zoning feature as discussed previously.

The plasma expands from the target surface in a planar manner until it reaches a distance equal to half the focal spot size of the laser pulse. Lateral expansion of the ablated material from the centre of the focal spot expands away at the sound speed. Once the material has expanded to the outermost region of the focal spot, a distance of half the focal spot size, the plasma begins to expand in a spherical manner. Therefore the experiment is well simulated by 1d simulations close to the target surface.

To study helium-like ion stages in detail a collisional-radiative model is used for hydrogen, helium and lithium-like ion stages where atomic data is generated in EHYBRID from a screened hydrogenic approximation. This provides detail on the helium-like populations by taking into account other important routes for electron population flow [66]. The remainder of the ion stages are calculated using the modified Griem’s model as it provides a more rapid calculation. This model treats the problem as a simple two-level system, a ground state and an
excited state with an energy corresponding to a fractional principle quantum number [76]. A Thomas-Fermi model is used to model the free electrons and an ideal gas approximation is used for the nuclear part of the equation of state. This is to ensure that the regions close to the target surface containing large spatial gradients due to the short duration of the laser pulse are simulated accurately.

5.4 Simulation of wave amplification

A post-processor written by the author is used to calculate parametric instability growth based on theory developed in chapter 2. The post-processor uses a $n_e$, $T_e$ and $L_n$ to determine the temporal evolution of the maximum convective amplification factor for each spatial location in the plasma. The post-processor calculates parametric amplification at $n_{cr}$ due to resonance absorption, and parametric growth at and below $n_{cr}/4$ due to TPD and SRS. This allows the calculation of the maximum electric fields associated with plasmons amplified above the thermal level. The amplified electric fields due to parametric processes at $n_{cr}$ and at and below $n_{cr}/4$ are limited by wavebreaking. Damping of the plasmons uses electron-ion collision rates for a warm plasma. This rate is a function of electron temperature and average ionisation $\bar{Z}$.

Firstly picosecond pulse is considered where there is an absence of a pre-formed plasma through which the laser must propagate. A high contrast $10^6$ pulse is used. Therefore it is assumed plasma is not pre-formed. The electric field strengths are compared with those generated by low contrast picosecond simulations. Double pulse simulations are post-processed in the same manner. Amplified plasmons are studied at the time the femtosecond pulse interacts with the extending plasma plume generated by the picosecond pulse for interpulse time delays of 50 and 200 ps. Plasmon damping and suppression is considered in each case.
5.5 Picosecond pulse + pedestal

5.5.1 Hydrodynamics

In this section hydrodynamic simulation results are presented for laser intensities of $3 \times 10^{16}$ W cm$^{-2}$, $3 \times 10^{15}$ W cm$^{-2}$ and $8 \times 10^{14}$ W cm$^{-2}$ picosecond pulses with a pedestal contrast of better than $10^6$ (high contrast) and $10^3$ (low contrast). Figure 5.1 presents the temporal evolution of the electron density, electron temperature and population density of the $n = 3$ state for various distances above the target surface for a high contrast picosecond pulse with an intensity of $3 \times 10^{16}$ W cm$^{-2}$. $t = 0$ corresponds to maximum laser intensity. At the target surface the electron density sharply increases to approximately $1 \times 10^{23}$ cm$^{-3}$ at $t = 0$. The electron density then drops rapidly to approximately $1 \times 10^{22}$ cm$^{-3}$ within about 10 ps and then gradually reduces to $1 \times 10^{21}$ cm$^{-3}$ at $t = 500$ ps. At distances above the target surface electron density values remain at a low level until the “bulk” high density plasma extends into the region of space.

The low density plasma plume is composed of high temperature electrons from the initial surface of the target. Electron temperatures at the target surface rapidly rise above 1 keV during maximum laser irradiation and then rapidly reduce to below 100 eV at $t = 10$ ps. At distances above the target surface the plasma corona remains at high temperatures above 1 keV until the bulk high density plasma extends into this location. Temperatures then quickly fall below 10 eV.

The temporal evolution of population density for the $n = 3$ state at the target surface, as presented in figure 5.1 d, shows that the maximum population density duration is very short, approximately 5 ps. For positions above the target surface the population density remains at high levels for the plasma corona. When the bulk high density colder plasma arrives the population density reduces signifi-
Figure 5.1: Temporal evolution of electron density (solid line), temperature (dashed line) and population density in the $n = 3$ level (dotted line) at [a), d]) 5 $\mu$m, [b), e]) 10 $\mu$m and [c), f]) 20 $\mu$m above the target surface.

Plasmon amplification takes place during laser irradiation and the level of amplification is dependent on plasma parameters and spatial gradients. Figure 5.2 shows a spatial profile comparison between simulations using a high contrast picosecond pulse and a pedestal only pulse for a focal spot size of 15 $\mu$m. Spatial profiles are taken at laser intensity maximum. Electron density and temperature gradients are significantly smaller for high contrast picosecond pulses when compared to pedestal only pulses. This is due hydrodynamic motion occurring significantly. Therefore for clean the picosecond pulses the highest population density occupies the low density coronal plasma.
Figure 5.2: a) $n_e$ and b) $T_e$ spatial profiles at maximum intensity for a high contrast picosecond pulse (dashed line - $I_L = 3 \times 10^{16}$ W cm$^{-2}$) and pedestal only pulse (solid line - $I_L = 3 \times 10^{13}$ W cm$^{-2}$) with a focal spot size of 15 $\mu$m.

over a longer time scale for pedestal pulse irradiation. The plasma extends to approximately 16 $\mu$m above the target surface for pedestal pulses whereas for high contrast picosecond pulses the plasma extends to approximately 4 $\mu$m above the target surface. The density gradients, electron temperatures and plasma extent all have a significant role in plasmon amplification.

Figure 5.3 presents temporal profiles of electron densities and temperatures for the three laser intensities for differing heights above the target surface. At $t = 0$ ps the electron densities sharply rise to maximum values for all three laser intensities for low contrast pulses. At times $t > 0$ ps the densities fall gradually. At 10 and 20 $\mu$m above the target surface as shown the electron densities rise to maximum values at times after $t = 0$ ps as plasma expands into these regions and remain uniform over 500 ps as shown in figure 5.3. The maximum electron density value is dependent on the height above the target surface, reducing at greater distances. Electron temperatures at the target surface rise to a maximum at $t = 0$ ps to values approaching 1 keV for the highest laser intensity. Temperatures sharply decrease at the target surface at times $t > 0$ ps where heat conduction is at a maximum due to corresponding high densities. At heights above the target temperatures reach a maximum as plasma expands into these regions in a similar manner to the electron densities. Temperatures gradually reduce after $t = 0$ ps.
with a gradient dependent on the thermal heat conduction and therefore electron density. For greater distances above the surfaces where electron densities are lower the gradient reduces. On comparison to clean picosecond pulse simulations the electron densities and temperatures follow different trends.

Figure 5.4 represents the $n = 3$ population densities for low contrast picosecond pulses. $n = 4$ population densities follow the same trends but the population fraction is lower. Population densities sharply peak at $t = 0$ ps at target surface and then rapidly decrease to low values within a few picoseconds. At distances above the target surface maximum values occur later in time dependent on the
expansion speed of the plasma. The length of time the \( n = 3 \) state is populated increases at greater distances where full width at half maximum becomes 100 ps at 20 \( \mu \)m. Comparing to clean picosecond pulses the \( n = 3 \) state is populated at times where the electron densities are greater and temperatures are lower for low contrast picosecond pulses. Maximum population densities are not found in the low density plasma plume for low contrast picosecond pulses. Maximum population densities occur in higher density plasma where \( n_e > 10^{17} \text{ cm}^{-3} \) for low contrast picosecond pulses.

Figure 5.5 shows spatial profiles of electron densities and temperatures during
Figure 5.5: Spatial profiles of the electron density and temperature at $t = -2$ ps (solid line), $t = 0$ ps (dashed line) and $t = 2$ ps (dotted line) for a low contrast pulse. Laser intensities are $3 \times 10^{16}$ W cm$^{-2}$ [a), d)], $3 \times 10^{15}$ W cm$^{-2}$ [b), e)] and $8 \times 10^{14}$ W cm$^{-2}$ [c), f)].
picosecond pulse irradiation for low contrast picosecond pulse simulations. The spatial distribution of electron density for all intensities changes very little during irradiation by the picosecond pulse, shown in figures 5.5a, b and c. Density gradients are particularly steep between the target surface to approximately the critical density and lower where gradients become shallower. For the highest laser intensity electron density occurs at target surface and falls rapidly from solid density and extends to approximately \(16 \mu m\). The size and mass of the target acts as a boundary to plasma expansion into the target. Steeper gradients are observed in the plume for lower laser intensities as the spatial extent of the plasma is reduced. The blow off plasma, created by the pedestal, containing electron densities less than or equal to quarter critical density occupies a significantly larger volume than plasma with densities approximately equal to critical density. The intense pulse interacts and loses energy to this extended plume region, via inverse bremsstrahlung and possibly non-linear processes, before dumping its energy at the critical surface or being reflected back into the lower density plasma.

Electron temperatures are shown in figures 5.5d, e and f. As the picosecond pulse arrives, the plasma increases in temperature significantly from pre-formed plasma values, creating very steep gradients between the target surface and approximately \(2 \mu m\) above it. At distances above \(2 - 3 \mu m\) the electron temperature becomes uniform until it rapidly decreases to values similar to pre-formed plasma temperatures. Only at times after the laser pulse does the temperature become completely uniform beyond the steep gradient. This occurs because the high temperature plasma heated by the picosecond pulse expands adiabatically into colder pre-formed plasma.

The spatial distribution of average ionisation \(\bar{Z}\) is an important value in calculating the electron-ion collision rate. It partially dictates the maximum convective amplification and the rate of plasmon damping. \(\bar{Z}\) is shown in figure 5.6 during picosecond pulse irradiation for the three laser intensities. For the highest laser
intensity $\bar{Z}$ increases rapidly at the target surface from zero (no ionisation) to $\bar{Z} = 11$ at the target surface indicating the majority of ions are He-like. Values remain uniform at $\bar{Z} = 11$ until approximately 8 $\mu$m above the target surface where $\bar{Z}$ gradually reduces to 8 where the majority of ions are B-like. The reduction of $\bar{Z}$ at the extent of the plasma is due to lower electron densities resulting in reduced collisional ionisation. For lower laser intensities $\bar{Z}$ reaches a maximum value of 11 during maximum laser irradiation. Values do not remain uniform at greater distances as electron densities as shown in figure 5.5 reduce significantly above target surface. $\bar{Z}$ reduces to values of 4 - 5 at the plasma extent.
5.5.2 Plasmon amplification

In this section the presence of plasmons amplified by parametric instabilities will be discussed. Using the spatially and temporally evolving values of \( n_e, T_e, L_n \) and \( \bar{Z} \) generated by EHYBRID the presence, amplification factors and spatial/temporal damping of plasmons are calculated. Thermal level plasmon electric fields without amplification are presented for low contrast picosecond pulses. Threshold parameters for TPD and SRS are discussed for high and low contrast picosecond pulses. The electric field strengths \( E_0 \) associated with instability and resonant absorption amplified plasmons is then presented for the three laser intensities. Plasmon decay and suppression is then discussed.

Figure 5.7 shows values of the thermal plasmon electric field strength at peak irradiance for all laser intensities. These values do not take into account any amplification due to non-linear processes. The values shown are dependent on electron temperature and the Debye length. For the highest laser intensity the greatest values of \( E_0 \) are found at the target surface with values exceeding to \( 10^9 \) V \( \text{cm}^{-1} \). At distances above target surface where electron densities \( n_e < n_{cr} \) \( E_0 \) reduces significantly to \( E_0 < 1 \times 10^7 \) V \( \text{cm}^{-1} \), reducing gradually to \( \approx 1 \times 10^5 \) V \( \text{cm}^{-1} \) at the plasma extremity. For lower laser intensities maximum \( E_0 \) occurs at target surface where electron densities and temperatures are at their highest. As with the highest intensity calculation, \( E_0 \) reduces to \( E_0 < 1 \times 10^7 \) V \( \text{cm}^{-1} \) at distances where \( n_e < n_{cr} \). \( E_0 \) then reduces to \( \approx 1 \times 10^5 \) V \( \text{cm}^{-1} \) with a gradient dependent on the plasma size and therefore the laser intensity.

The threshold parameters for TPD and SRS instabilities are calculated using the EHYBRID simulations. SRS is not present in all simulations due to it’s higher threshold 2.7.2. Figure 5.8 shows the spatial distribution the threshold parameter \( \xi = (1/12)(v_{os}/v_{te})^2 k_0 L_n \) for TPD from chapter 2, at times during picosecond pulse irradiation. This value must be equal to or greater than unity for significant TPD to occur. For low contrast pulses the TPD threshold is
Figure 5.7: Spatial distribution of thermal plasmon electric field strengths at $t = -2$ ps (solid line), $t = 0$ ps (dashed line) and $t = 2$ ps (dotted line) for a low contrast pulse.

exceeded for the highest laser intensity for all times and for all distances where $n_e \leq n_{cr}/4$ during picosecond pulse irradiation. For the medium laser intensity the threshold is met at $t = -2$ ps. However for the lowest laser intensity the TPD threshold is not exceeded for any time during the picosecond irradiation.

Once the two-plasmon decay threshold has been exceeded then the convective amplification threshold can be used to find the maximum amplified electric field of the plasmons. Figure 5.9 shows the amplified plasmon electric field strengths due to two-plasmon decay for all laser intensities. The fields are limited by wave breaking. For the lowest laser intensity there is no amplification as two-plasmon decay does not occur. For the highest laser intensity the thermal plasmon field increases substantially from thermal levels to approximately $5 \times 10^8$ V cm$^{-1}$. For the medium laser intensity plasmon electric field strength are amplified to values of approximately $5 \times 10^8$ V cm$^{-1}$ at electron densities below quarter critical density, significantly lower than values inferred from data. This, along with lower laser intensities of $8 \times 10^{14}$ W cm$^{-2}$, suggests that the electric fields as inferred from data cannot be driven by this mechanism.

However in all cases the laser power does exceed the self-focussing threshold and therefore it is probable self-focussing occurs [77]. In hydrodynamic simu-
Figure 5.8: Spatial distribution of the two-plasmon decay threshold at $t = -2$ ps (solid line), $t = 0$ ps (dashed line) and $t = 2$ ps (dotted line) for a low contrast pulse. Laser intensities are a) $3 \times 10^{16}$ W cm$^{-2}$, b) $3 \times 10^{15}$ W cm$^{-2}$ and c) $8 \times 10^{14}$ W cm$^{-2}$. Note the different scales.
Figure 5.9: Spatial distribution of the amplified plasmon electric field strengths due to two-plasmon decay at $t = -2$ ps (solid line), $t = 0$ ps (dashed line) and $t = 2$ ps (dotted line) for a low contrast pulse. Laser intensities are a) $3 \times 10^{16}$ W cm$^{-2}$, b) $3 \times 10^{15}$ W cm$^{-2}$ and c) $8 \times 10^{14}$ W cm$^{-2}$ respectively.

The growth rate of self-focussing is around $0.1$ ps$^{-1}$ whereas, in contrast, the growth rate of two-plasmon decay is $11$ ps$^{-1}$. This suggests that TPD is quickly established and plasmon amplification occurs throughout the laser pulse whereas self-focussing occurs on a longer timescale. For spot size to filament width ratios of 2.5 and 5.0 for laser intensities of $3 \times 10^{15}$ W cm$^{-2}$ and $8 \times 10^{14}$ W cm$^{-2}$ respectively the two-plasmon decay threshold is exceeded and plasmons are amplified to $1 \times 10^8$ V cm$^{-1}$.

In addition resonant absorption at the critical surface drives intense plasma
waves as discussed in 2.7.1. The peak fields are shown in table 5.1. Any spectroscopic observation of an intense plasma wave relies on the spatial and temporal coincidence of the plasmons with intense spectral line emission. The data available shows modulations on the He $\beta$ ($n = 3 \rightarrow 1$) and the He $\gamma$ ($n = 4 \rightarrow 1$) transitions. This is discussed in chapter 6.

### 5.5.3 Plasmon damping/suppression

Figure 5.10 a) shows the spatial distribution of plasmon collisional damping rates $\Gamma_C$ for all laser intensities. The calculations are used to estimate plasmon damping both spatially and temporally. As the electron densities and ionisation levels increase further from the target surface $\Gamma_C$ reduces significantly. Indeed the group velocity of the plasmons $v_g$ is such that the amplified plasmons can propagate through the low density plasma without significant damping.

The spatial locations above the critical density and the quarter critical surfaces have approximately uniform values of $\Gamma_C < 0.5$. Figures 5.10 a and b show the temporal and spatial decay of plasmons from an amplitude value of $5 \times 10^8$ V cm$^{-1}$ where $\Gamma_C = 0.5$. The spatial decay is a function of the collisional damping rate and the group velocity of the plasmons. Close to the target surface the group velocity can reach values of up to $\approx 0.2c$. Therefore for even though the plasmons may decay quickly in time compared to space the high amplitude plasmons reach the extremity of the plasma expansion and are Landau damped significantly. This occurs within approximately 2 ps, before $n = 3$ ion population

<table>
<thead>
<tr>
<th>Laser intensity (W cm$^{-2}$)</th>
<th>Maximum electric field (V cm$^{-1}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$3 \times 10^{16}$</td>
<td>$6 \times 10^8$</td>
</tr>
<tr>
<td>$3 \times 10^{15}$</td>
<td>$2 \times 10^8$</td>
</tr>
<tr>
<td>$8 \times 10^{14}$</td>
<td>$6 \times 10^7$</td>
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Table 5.1: Maximum electric field amplification due to resonance absorption.
density maximum. At times up to approximately 25 to 30 ps the electric fields values remain above $5 \times 10^7$ V cm$^{-1}$ in the convective frame. The population density of the $n = 3$ and 4 levels occur later in time at greater distances above the target surface. The plasmon amplitudes in the convective frame will decrease as the plasma expands due to collisional damping. Therefore at greater distances above the target surface the electric field imposed upon the ions will reduce.

### 5.6 Summary

In this chapter hydrodynamic simulations from EHYBRID using experimental laser and target parameters were presented for picosecond plus pedestal. Plasmon amplification due to parametric instabilities and subsequent decay by collisional damping was calculated. High contrast picosecond pulse plasmon amplification does not occur due to the absence of a pre-formed plasma. For low contrast picosecond pulse and double pulse simulations electric field strengths are found to be greater than $5 \times 10^8$ V cm$^{-1}$, limited by large amplitude wave breaking. Electric field strengths inferred from data are approximately $1(\pm 0.1) \times 10^8$ V cm$^{-1}$. The plasmon electric field strengths reduce above target surface due to
collisional damping. Maximum population in the $n = 3$ and 4 levels coincides with maximum electric field strengths associated with the plasmons. This suggests an interaction between the 1s3p and 1s4d configurations with the plasmon field close to the target surface. This supports the spatial locations of the modulations as found in data.
Chapter 6

Lineshape simulations

6.1 Introduction

In this chapter the atomic kinetics code CRETIN [78] is described and used to calculate line shapes using simulated hydrodynamic data as presented in the previous section. There is a need to test the confidence of the hydrodynamic simulations. The approach is to use cross sections taken at or near the target surface where pressure broadening dominates and where the L-dips are observed. Figure 6.1 presents a flowchart of this chapter. Firstly the CRETIN code is used to calculate the K-shell line emission as a function of time and electron density. Simulated time integrated K-shell line shapes are then compared to data. Secondly a post-processor (developed by the author) is used to calculate the L-dip positions and structures for different plasmon frequencies, electric field strengths and number of quanta involved in the interaction. Composite spectra are formed using the post-processor output and line shapes calculated using single $n_e$ and $T_e$ and compared to data.
6.2 CRETIN

The atomic kinetics code CRETIN was initially developed for use in astrophysical modelling. The code has also been used extensively to model a number of experiments in areas such as inertial confinement fusion, magnetic fusion and laser-produced plasmas. It is able to solve atomic level populations, kinetics and radiation distribution self consistently in a multidimensional non-local thermodynamic equilibrium (NLTE) plasma. CRETIN calculates most common atomic processes where forward rates, such as excitation and ionisation, are calculated from information included in an atomic model or tabular input. Inverse processes are obtained from detailed balance relations. Collisional excitation and ionisation rates are calculated from fits to Maxwellian-averaged strengths as a function of electron temperature. The equations of radiative transfer (treated separately in three “phases”; continuum, bound-bound transitions and spectral construction) and statistical equilibrium are linearised and simultaneously solved by iteration.

The linearised equations can be calculated in 1, 2 or 3 dimensions using a cartesian, cylindrical or spherical coordinate system depending on the needs of
the problem in question. The calculations are performed using a Lagrangian mesh if the plasma is allowed to evolve in time for example when supplying hydrodynamic data. Due to the Lagrangian configuration, where each cell must contain the same mass throughout the simulation, cells flow past the line of sight in the experimental frame of reference. To calculate lineshapes for a particular location in the experiment frame of reference using a time-dependent evolution the simulation must be completed and then be remapped onto a Eulerian grid. The atomic data used can either be calculated by CRETIN using a screened hydrogenic approximation or supplied by a data file. The HULLAC [79] code can be used to calculate accurate atomic data for CRETIN. The HULLAC data contains multiple transitions including the neighbouring Li-like satellites. Output such as the temporal evolution of line strengths, ratio of magnitudes between neighbouring lines and spectral broadening is similar when using screened hydrogenic approximations and the HULLAC data for a given hydrodynamic input [66].

6.2.1 1-dimensional time-dependent atomic kinetic calculations

In this section the results of CRETIN simulations driven by hydrodynamic data calculated using the EHYBRID code in chapter 5 for a single focal spot size of 15 µm are presented. The purpose of this is to calculate the amount of photons/mode (the brightness per unit frequency) emitted as a function of time and to simulate the time integrated line shapes. The He β and γ transition is studied for a number of locations close to the target surface to compare with data and the plasmon amplification results. In this case the plasma is represented in 1-dimension extending parallel to the angle of observation i.e. perpendicular to plasma expansion. The data provided to CRETIN consists of time varying electron density, electron temperature and ion temperature for positions up to 20 µm above the target surface. The plasma size is fixed to 15 µm for all positions.
Figure 6.2: Temporal evolution of the emitted photons/mode and associated electron densities for a spatial location of a) 5 µm, b) 10 µm and c) 20 µm. Solid lines represent \( n_e \), dashed lines He \( \beta \) transition, and dotted lines represent the He \( \gamma \) transition. Note the photons/mode scale changes for each graphic.

Therefore the amount of photon reabsorption due to the plasma extent is equal for all three focal spot sizes. The simulations are run using a screened hydrogenic approximation to reduce the time for each CRETIN iteration. CRETIN is run in time-dependent non-local thermodynamic equilibrium conditions (NLTE).

Figure 6.2 presents the emitted photons/mode for the He \( \beta \) and \( \gamma \) transitions and associated electron densities as a function of time for a 15 µm spot size picosecond only pulse interaction for spatial locations of 5 µm, 10 µm and 20 µm above the target surface. The maximum emission for all spatial location
occurs after $t = 0$ where the hot ionised emitting plasma enters this region of space. At 5 $\mu$m above the target surface He $\beta$ and $\gamma$ emission has two maxima. The first occurs approximately 4 ps after maximum picosecond pulse irradiation. At this time the He $\beta$ and $\gamma$ lines are being emitted from electron densities of $n_e \simeq n_{cr}/4$. This is during the period where amplified plasmons occur. The secondary emission maximum occurs approximately 30 ps after the picosecond pulse. The emission is strongest during this period. The He $\beta$ and $\gamma$ lines are emitted from higher density plasma where $n_e \simeq 6 \times 10^{21}$ cm$^{-3}$. The total He $\beta$ and $\gamma$ emission time at this height is approximately 60 ps. For greater heights above the target surface there is a similar trend. However at 20 $\mu$m the magnitude of two emission maxima are similar. The emission reduces in magnitude at greater distances above the target surfaces and increases in duration.

### 6.2.2 1-dimensional time-dependent integrated lineshapes

In this section the time-integrated line shapes calculated by CRETIN using the output of EHYBRID are presented. The spectra are a sum of the entire time history of the plasma emission at specific distances above the target surface. Figure 6.3 shows a comparison between calculated time-integrated He $\beta$ line shapes and data at different heights above the target surface for a picosecond only pulse. The laser focal spot size is 15 $\mu$m. The calculated spectra are formed using the time-dependent line strengths presented in section 6.2.1. This includes emission from low electron densities where $n_e \simeq n_{cr}/4$ and high electron densities ($n_e > 10^{21}$). The calculated line shapes for each height are a good match to data. Close to the target surface the calculated He $\beta$ line shapes are indented at the line centre. This is also found in data. As emission from the rear of the plasma (i.e. the portion of the plasma furthest from the viewer) travels through the plasma toward the viewer the lineshape becomes more “flat-topped” due to absorption. Greater contributions to the lineshape occur for larger distances between the viewer and
the emitting plasma element. The sum of the emitting elements over the extent of the plasma results in a lower level of emission at the line centre than in the wings.

Figure 6.3 shows that the calculated time-integrated K-shell line profiles are similar to those found in data at heights close to the target surface for picosecond only pulses and double pulses. This confirms that the hydrodynamic data output from EHYBRID is a good approximation to the spatial and temporal evolution of the experimental plasma.
6.3 L-dip simulations

It is possible to reconstruct the L-dip structures as seen in data by post-processing the simulated lineshapes using the Intra-Stark spectroscopy framework as discussed in 3.5.1. As previously discussed the positions and structure of the L-dips within the spectra are dependent on the upper quantum number of the transition, the electron density from which the bound-bound line is emitted and the electric field strength associated with the resonant plasmons. The author has constructed such a post-processor to evaluate the magnitude and extent of L-dips for various electron densities (broadband plasmon frequency), electric field strengths and number of interacting quanta. For multi-quantum interactions of \( p = n_p \) quanta the L-dips are calculated for \( p = 1..n_p \) ensembles of radiators. Therefore second order L-dips \( (p = 2) \) overlap with first order L-dips \( (p = 1) \) in cases of low electron densities \( n_e \approx 10^{20} \text{ cm}^{-3} \). The electron densities can be chosen by the user to be either single or multi valued. The electric field strengths, which govern the half-width and depth of the dips where \( \Delta S/S \approx 1 \), can also be treated to be single or multi valued. The purpose of this is to analyse the temporal evolution of the plasma properties and the effect of plasmon damping in dynamic resonance on L-dip observation.

6.3.1 Single \( n_e, T_e \) and \( E_0 \) simulations

L-dips in this section are created using the single valued electron densities, electric field strengths and number of interacting quanta as inferred from data shown in tables 4.5 and 4.6. These calculated line shapes are post-processed to include L-dips. The line shapes are calculated using single \( n_e \) and \( T_e \) for a 0-dimensional plasma with no optical depth. Figure 6.4 presents an example of a comparison between He \( \beta \) line shapes calculated using single \( n_e \) and \( T_e \) and data using a 15 \( \mu \text{m} \) laser focal spot size. Cross sections are taken at target surface for picosecond
Figure 6.4: Calculated single \( n_e, T_e \) He \( \beta \) line shapes (dashed red line) compared to data (solid black line) for picosecond only pulses and \( \Delta t = 200 \) ps double pulse.

only data and double pulse data with different time delays. In this case \( T_e \) is chosen to be 800 eV. The letters “ps” denote the picosecond pulse and “fs” the femtosecond pulse with a time delay of “(\( \Delta t \))”. Satellites found on the low energy wing of the line are not modelled. This results in a discrepancy between data and calculation.

Figure 6.5 presents a comparison of He \( \beta \) data and simulated L-dips for a laser focal spot size of 15 \( \mu \)m for single and double pulse. Cross sections are taken at 6 \( \mu \)m above the target surface. The post-processor alters the single \( n_e, T_e \) line shapes to include L-dips. The composite spectra have the same underlying shape
as the CRETIN simulations. The L-dips do not alter the line shape significantly. The overlayed simulated post-processed lineshape is particularly similar to data. The majority of L-dips and the structures they form due to the constructive and destructive interaction between one another are accounted for. In the majority of data cross sections close to the target surface L-dips are suppressed in data close to the line centre where the plasma optical depth causes photon reabsorption. For an optically thin plasma the emitted lineshape will not be affected by reabsorption and a Lorentzian lineshape will be observed. Furthermore L-dips will not be suppressed for optically thin plasma. As the single $n_e$, $T_e$ CRETIN simulations are run in zero-dimensions the optical depth is effectively zero and therefore no reabsorption takes place.

As a comparison figure 6.6 presents a comparison of He $\beta$ data and simulated L-dips for a laser focal spot size of 50 $\mu$m for single and double laser pulses. Similarly to the simulated L-dips for focal spot sizes of 15 $\mu$m as shown in figure 6.5 the simulations account for the majority of L-dips and structures they create in data. Similarly figure 6.7 shows a comparison between L-dips and data for He $\gamma$ cross sections. A good agreement is found for all laser focal spot sizes. In this case the electron densities inferred from cross sections are comparable to those found from He $\beta$ cross sections. However the number of quanta involved in the interaction is reduced and therefore the complexity of the spectrum is reduced. This is most probably due to a lower density of He $\gamma$ radiators compared to He $\beta$ as shown in figures 6.2. The fact that no L-dips are observed in He $\gamma$ lineshapes for double pulse data is most likely due to the “washing out” of the L-dips produced by dynamic resonance with amplified plasmons created by the primary pulse. Due to the lower He $\gamma$ emission and fewer plasmon quanta involved in resonance compared to He $\beta$ the emission due to the secondary pulse masks the L-dips in time integration. Furthermore the onset of SRS will create a broadband frequency amplification of electron plasma wave modes. This will affect the single frequency TPD amplification at $n_e = n_{cr}/4$. 

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Figure 6.5: Comparison of simulated L-dips and 15 μm focal spot He β data for cross sections taken at 6 μm above the target surface for a) primary pulse only and double pulses with b) Δt = 50 ps and c) Δt = 200 ps. Solid lines and dashed lines indicate data and simulation respectively.
Figure 6.6: Comparison of simulated L-dips and 50 µm focal spot He β data for cross sections taken at 6 µm above the target surface for a) primary pulse only and double pulses with b) Δt = 50 ps, c) Δt = 200 ps and d) Δt = 400 ps. Solid lines and dashed lines indicate data and simulation respectively.
Figure 6.7: Comparison of simulated L-dips and single pulse He $\gamma$ data cross sections taken at 7 $\mu$m above the target surface for a spot size of a) 15 $\mu$m, b) 50 $\mu$m and 100 $\mu$m. Solid lines and dashed lines indicate data and simulation respectively.
6.3 Lineshape simulations

Figure 6.8: The effect of collisional damping during and after laser irradiation. Individual contributions to the spectrum a) are calculated at 0.0 ps (solid black), 2.5 ps (dashed red) and 5.0 ps (dotted blue) after laser irradiation. A time integrated spectrum b) is calculated using all the contributions.

6.3.2 Effects of collisional damping

The competing effect of laser driven instabilities and collisional damping limits the presence of electron plasma waves to a narrow time window centred around the laser peak intensity. Composite lineshape calculations presented in figure 6.8 illustrate how the Intra-Stark structures (L-dips) are only present when the plasma waves are strong. Figure 6.8 a shows composite spectra at times during and after peak laser irradiance and figure 6.8 b shows a time integrated spectrum. A single frequency plasma wave is amplified during laser irradiation. From simulations $E_0 \approx 2 \times 10^8$ V cm$^{-1}$ during laser irradiation. This value decreases to approximately $1 \times 10^8$ V cm$^{-1}$ 2.5 ps after maximum irradiation and $4 \times 10^7$ V
6.4 Lineshape simulations

\[ \text{cm}^{-1} \] after 5.0 ps. This results in an Intra-Stark spectrum generated by maximum amplitude plasma waves. For the L-dips to be observable the He-like emission must be short, on the order of the laser pulse. This is the case for picosecond only pulses, not for double pulses.

6.4 Summary

In this chapter the temporal evolution of resonance line strength and time integrated K-shell spectra were simulated by the CRETIN code. Simulations were driven using EHYBRID output. For picosecond only pulses the He $\beta$ and He $\gamma$ lines are emitted from electron densities $n_e \leq n_{cr}/4$ close to the target surface during irradiation. The resonance lines also emit from higher electron densities at later times. Simulated time integrated spectra are compared to data. Simulations match to data cross sections taken close to the target surface for the three laser intensities and interpulse time delays. This supports the EHYBRID hydrodynamic simulations. Simulated spectra are post-processed to include L-dips. Parameters inferred from data are used to construct L-dips structures and positions. The composite spectra are compared to data. The simulated L-dip structures are found to match those found in data.
Chapter 7

Conclusions

In this thesis a study and investigation of high field effects on laser produced plasma spectral line shapes was presented. This includes the effects of electron plasma waves on bound-bound transitions. This chapter provides a summary and conclusion to the experimental and computational work presented in the thesis. Possibilities for further work are given.

7.1 Conclusions

With the experimental design and diagnostic setup presented in chapter 4 high field effects were successfully observed. These effects are clearly seen in the time-integrated He $\beta$ transition and suggested by the He $\gamma$ data. These are the first observations in the He-like series. The He $\beta$ was chosen as the transition is positioned in a relatively clean part of the K-shell spectrum and the associated dielectronic satellites are well documented. The difference between the laser wavelength and transition wavelength demanded high spectral resolution and dispersion. This with the need for spatial resolution required the use of a high performance toroidal spectrometer.
The structure of the modulations consists of a number of regular intensity maxima and minima on the spectral line. The modulations are observed in data close to the target surface for picosecond laser pulses. At greater distances above the surface they reduce in number and amplitude. When a femtosecond pulse is introduced after the picosecond pulse the modulations are not as strong.

Data analysis shows that these features are not due to known dielectric satellite positions and they are not due to noise or structure in the detector. This suggests that the features were due to some property of the plasma. The regularity of the structures/modulations suggests an interaction between electron plasma waves and Stark split energy levels. The data taken closest to target surface shows the most structure and this structure was richest at higher intensities.

The modulations are analysed using the Intra-Stark framework as developed by Oks. Using this theory the frequencies and electric field strengths associated with the electron plasma waves are directly inferred from the modulation location and structure. This interpretation suggests that the electron densities at which the Intra-Stark processes occur are close to and just below quarter critical density. The inferred electric field strengths are approximately $10^9$ V cm$^{-1}$. Data interpretation indicates that the electric field strengths are at a maximum close to the target surface and at highest laser intensity. The number of modulated features were matched by theory only if theory was extended to include multiple quanta (up to 4) interacting with the Stark sublevels.

EHYBRID hydrodynamic simulations of the experiment were used to estimate the electron density and temperatures in the experiment. Post-processing the simulations to calculate the presence of laser driven plasma waves shows that plasma waves were present. These simulations give an indication of the likely amplification and damping of these plasma waves as a function of space and time. For the highest laser intensity only, the plasma waves were driven to saturation, and the calculated fields are similar to those inferred from the experimental re-
results. At lower intensity, intensity enhancement is necessary to achieve fields approaching those inferred. The simulations show that the plasma length scales and temperatures were suitable for self focusing. If this mechanism occurs then the laser intensities are sufficient to drive plasma waves, and calculated fields match those inferred from measurement.

The simulations show that TPD processes are efficiently amplified and that this is due to the long length-scale plasma created by the low contrast in the laser pulse. However, it is unlikely that SRS occurred due to its higher threshold for picosecond laser pulses. However for femtosecond pulses the SRS mechanism is more likely to occur and a broadband frequency of electron plasma modes will be amplified. This is the most likely reason why L-dips reduce in number and magnitude in double pulse data. Damping of the electron plasma waves by electron-ion collisions is significant, effectively quenching the electric field after laser irradiation. It is likely that the electric fields inferred from data are driven at the time of picosecond laser irradiation by TPD and associated electron plasma waves.

The high electric field strengths coincide with the population of the $n = 3$ and $n = 4$ levels for the low contrast laser pulse. CRETIN simulations of the temporal evolution of the H-like emission suggests that the He $\beta$ and He $\gamma$ lines are emitted from electron densities of approximately quarter critical density during and soon after picosecond pulse irradiation. Time integrated lineshape simulations are found to match data for different spatial locations, laser intensities and time delays. Intra-Stark modulations (L-dips) were calculated using the parameters inferred from data and superimposed onto the simulated lineshapes to form composite spectra. The simulated structures formed by the overlapping L-dips reproduce those found in data.
7.2 Further work

To characterise the plasma fully it would be necessary to use additional complementary diagnostics. Pinhole cameras should be used to measure the shape of the expanding plasma plume. In addition the parametric instabilities could be examined using temporally resolved Thomson scattering. Scattering measurements could be used to support simulations of parametric instabilities. With regard to the crystal spectrometers additional resonance line emission particularly of the H-like resonance lines would further characterise the plasma. The higher electric field susceptibility of H-like transitions would allow L-dip structures to be formed for low electric field strengths. The observation of L-dips created by high electric field strengths in H-like systems would support the theory of dynamic resonance in He-like systems. Furthermore electric field strengths and electron plasma wave frequencies measurements can be inferred with greater accuracy due to the reduced complexity of the H-like Intra-Stark spectra. Furthermore by introducing polarisers into the spectrometer system it would be possible to mutually exclude the $\pi$ and $\sigma$ portions of the L-dip spectra. This would reduce the complexity of the spectrum and allow for increased accuracy. This can be achieved by placing a second crystal between the HDTS and the CCD detector. The crystal Bragg reflects only the $\sigma$ polarised emission onto a secondary CCD detector and allows unpolarised light to pass through. The $\sigma$ polarised spectrum can be directly compared to the unpolarised spectrum from a single shot. For example the Al He $\beta$ line a 10 $\mu$m thick mica crystal ($2d \approx 19.84$ Å, 002) positioned with an angle of 42° will reflect second order $\sigma$ polarised emission approximately 70 times stronger than $\pi$ polarised emission. The thickness of the crystal is an issue as mica strongly absorbs at this wavelength. For a crystal thickness of 10 $\mu$m approximately 75 % of the unpolarised emission is absorbed [80].

In the ASTRA experiment the energy on target was fairly low and therefore the plasma emission was limited. Shot-integration was employed to increase the
signal. Using the ASTRA GEMINI upgrade it is possible to deliver short pulse high energy pulses onto target. Single shot data would be recorded reducing the complexity of the experiment. Furthermore using higher intensity beams would cause SRS to occur. This would generate a larger amount of amplified electron plasma waves and provide an extra diagnostic due to the scattered photons.

It would be particularly interesting to simulate and study the populations and lifetimes of the Stark sublevel states. By introducing this simulation parameter into the line shape model, L-dip spectra could be calculated with greater accuracy.

This diagnostic can be used to spatially locate parametric instabilities within short pulse laser produced plasma. It can be employed in experiments such as fast ignition ICF and wake field particle acceleration to diagnose the high frequency electric fields associated with laser driven electron plasma waves. In ICF and wake field experiments the electron densities of the preformed plasma or gas jet target are approximately $10^{18} - 10^{20}$ cm$^{-3}$ with large density scale lengths and high ionisations. In these conditions the threshold for the onset of parametric instabilities is low. Raman scattering is commonly used to diagnose the plasma. The spectroscopic determination of parametric instabilities and measurement of high frequency fields would be a useful complimentary diagnostic.
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