CHARACTERIZATION OF PLASMAS PRODUCED BY PICOSECOND AND NANOSECOND LASER PULSES

by

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Abstract

This thesis reports on an experimental study of high power laser interaction with matter. The interaction physics with long scalelength plasmas was studied in a nanosecond regime relevant to Inertial Confinement Fusion (ICF). The interaction of picosecond laser pulses with micron-sized plasmas was investigated in a high intensity regime. Particular attention was devoted to plasma production and characterization using newly developed diagnostic techniques including time-resolved X-ray imaging and picosecond resolution X-ray spectroscopy.

An overview on laser-driven plasma instabilities is given with emphasis on stimulated Brillouin scattering, two plasmon decay and filamentation instability. Mechanisms for absorption and emission of electromagnetic radiation in inhomogeneous plasmas are described analytically. Non-linear effects in energy absorption mechanisms as well as dense plasma effects are taken into account. Numerical codes are employed to simulate plasma hydrodynamics and atomic physics. A detailed description of these codes is given along with a discussion of the limits of applicability to the experimental conditions under investigation.

Thin foil targets were used to generate long scalelength plasmas which were fully characterized in terms of electron density, using interferometric measurements, and electron temperature using time-resolved X-ray spectroscopy. Novel X-ray imaging techniques are described which allowed the first time-resolved X-ray images of the plasma to be acquired with a temporal resolution of ≈100 ps. Advanced numerical techniques were implemented and used in the analysis of interferograms to obtain fully three dimensional density profiles of the plasma. The suitability of these plasmas, for interaction studies relevant to ICF is discussed and experimental evidence of the onset of filamentation and whole-beam self-focusing even in the presence of laser beam smoothing is presented.

Short scalelength plasmas were produced irradiating solid targets as well as layered targets with high power, prepulse-free picosecond laser pulses. Detailed absorption measurements are presented which indicate the role of resonance absorption mechanisms in the laser-target interaction. These results are related to X-UV spectroscopic measurements with picosecond temporal resolution which enabled the first observation of transient heat propagation in solids in a picosecond regime. On the basis of these measurements, the interplay between absorption processes and electron heat transport is discussed in detail using analytical models and numerical simulation.
Acknowledgements and Role of the Author

The experimental work presented in this thesis was only possible due to the invaluable contribution of many people. First of all I would like to thank Prof. Oswald Willi for offering me the opportunity to join his research group and for constantly and enthusiastically suggesting a comprehensive selection of far-reaching ideas which, if I had taken note of them, would keep me engaged for the rest of my scientific career, if any, and beyond.

It is a great pleasure to thank the staff of the Central Laser Facility at the Rutherford Appleton Laboratory, where most of the experimental activity related to the thesis was carried out, for their enthusiastic and invaluable support. In particular, my thanks go to Colin Danson, Erol Harvey, Jim Lister, Dave Pepler, Dave Rodkiss and Ray Wyatt for their precious contribution and also for being understanding beyond any reasonable limit. I’m also deeply grateful to Marcos Protopapas for his invaluable contribution to the edition of the thesis.

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Role of the author

All of the experimental data presented in this thesis was obtained by the author during several experiments performed at the Rutherford Laboratory. The only exceptions are the interferograms shown in Figs.4.2.1-2 and Figs.4.6.1-2, the time resolved image shown in Fig.4.6.3 and the Brillouin scattering measurements of Fig.4.6.4. The analysis of the experimental data was entirely carried out by the author, with the following exceptions: thanks are due to Dr. Danilo Giuliani for his contribution to the analysis of the X-ray spectra of Fig.4.5.4, and to Dr. Antonio Giuliani and co-workers, for the study concerning the time resolved image of Fig.4.6.3. Finally, the primary role played by Dr. Tourage Afshar-rad, in the analysis of the time-resolved X-ray image of Fig.4.6.5 is acknowledged.
List of Symbols and Definitions

Frequently used symbols

- \( v \) : velocity \( \text{cm/sec} \)
- \( v_{th} \) : electron thermal velocity \( \text{cm/sec} \)
- \( v_s \) : ion sound velocity \( \text{cm/sec} \)
- \( v_q \) : electron quiver velocity \( \text{cm/sec} \)
- \( n, n_e, n_i \) : number density \( \text{cm}^{-3} \)
- \( n_{cr} \) : critical number density \( \text{cm}^{-3} \)
- \( T, T_e, T_i \) : temperature \( \text{deg}(K) \)
- \( \Theta, \Theta_e, \Theta_i \) : temperature in energy units \( \text{eV} \)
- \( k \) : wavenumber \( \text{cm}^{-1} \)
- \( \lambda \) : wavelength of e.m. wave \( \text{cm} \)
- \( \omega \) : angular frequency \( \text{rad/sec} \)
- \( \omega_p \) : plasma angular frequency \( \text{rad/sec} \)
- \( \nu_{ei} \) : e-i collision frequency \( \text{sec}^{-1} \)
- \( \nu_{ee} \) : e-e collision frequency \( \text{sec}^{-1} \)
- \( \mu_r \) : plasma refractive index \( 1 \)
- \( \varepsilon \) : dielectric permittivity \( 1 \)
- \( Z \) : charge state \( 1 \)
- \( A \) : atomic weight \( 1 \)
- \( m_i \) : ion mass \( g \)
- \( \tau_L \) : Landau damping constant \( \text{sec} \)
- \( \tau_{col} \) : collisional damping constant \( \text{sec} \)
- \( L \) : electron density scalelength \( \text{cm} \)
- \( L_v \) : fluid velocity scalelength \( \text{cm} \)
- \( \kappa_{ib} \) : inverse bremsstrahlung coefficient \( \text{cm}^{-1} \)
- \( F_e \) : electron thermal flux \( \text{erg/sec/cm}^2 \)
- \( F_e^{\text{max}} \) : free streaming electron thermal flux \( \text{erg/sec/cm}^2 \)
- \( \kappa_{SH}, \kappa_{FP} \) : thermal electron conductivity \( \text{cm}^{-1}\text{sec}^{-1} \)
- \( \kappa_{LD} \) : Landau damping conductivity \( \text{cm}^{-1}\text{sec}^{-1} \)
- \( \chi_0 \) : ionization potential \( \text{erg} \)
- \( A_{\text{u,l}} \) : Einstein spontaneous emission coefficient \( \text{sec}^{-1} \)
- \( \lambda_D \) : Debye screening length \( \text{cm} \)
- \( \eta \) : E-field strength parameter \( 1 \)
- \( \Delta \varphi, \delta \varphi \) : phase shift \( \text{rad} \)
Fundamental constants

\[ m_e = 9.1094 \times 10^{-28} \text{ g} \]
\[ m_p = 1.6726 \times 10^{-24} \text{ g} \]
\[ K_B = 1.3807 \times 10^{-16} \text{ erg/deg}(K) \]
\[ c = 2.99789 \times 10^{10} \text{ cm/sec} \]
\[ e = 4.8032 \times 10^{-10} \text{ statcoul} \]

Acronyms

ASE Amplified Spontaneous Emission
CG Cylindrical Geometry
EOS Equation of State
FD Fermi-Dirac
FFT Fast Fourier Transform
FI Filamentation Instability
ICF Inertial Confinement Fusion
ISI Induced Spatial Incoherence
LFG Line Focus Geometry
MCP Micro Channel Plate
PHC Pin Hole Camera
RPP Random Phase Plate
SBS Stimulated Brillouin Scattering
SF Self-focusing instability
SRS Stimulated Raman Scattering
SSD Smoothing by Spectral Dispersion
TF Thomas-Fermi
TPD Two Plasmon Decay
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Introduction

The interaction of intense laser light with matter provides a unique way to generate plasmas characterized by conditions of density and temperature otherwise impossible to achieve on a laboratory scale. These circumstances have given rise to a wide range of applications including, for example, inertial confinement fusion (ICF) and the X-ray laser. Besides the possible applications, however, the research on intense laser-matter interaction is now widely recognized as a novel and powerful tool to study the physics of atoms in conditions of extremely high electric field, never achieved before. In fact, the present day laser technology can provide laser pulses as short as $\approx 10^{-13}$ s with power as high as $\approx 10^{12}$ W and characterized by very high focusability which allows intensities of up to $\approx 10^{20}$ W/cm$^2$ and electric fields in excess of $\approx 10^{11}$ V/m to be achieved.

During the interaction of intense laser light with solid matter, a plasma is produced whose basic macroscopic features are mainly determined by the pulse duration and the light intensity and wavelength. Laser energy is then absorbed via a number of different mechanisms, the most important ones being collisional absorption and resonance absorption. Analytical models have been developed in the past years that allow the hydrodynamic expansion of laser-produced plasmas to be described. Numerical codes are presently in use world-wide which can provide simulations of the plasma temporal evolution from both an hydrodynamic as well as an atomic physics viewpoint. Only recently however, more attention has been devoted to a systematic experimental characterization of these plasmas which could provide, on one side, a satisfactory knowledge of plasma conditions for the analysis of interaction experiments, on the other side, an adequate test-bed for analytical and numerical models of plasma hydrodynamics and atomic physics.

The presently suggested schemes for direct drive ICF studies require laser drivers of typically 1 to 10 ns in duration. In the millimetre sized plasma produced in these conditions, collisional absorption is expected to be predominant over resonance absorption, resulting in most of the laser energy being absorbed in the low density corona surrounding the ICF spherical capsule. The interplay between the various mechanisms leading to absorption of electromagnetic energy in inhomogeneous plasmas will be described in detail. Non-linear effects in energy absorption will also be taken into account in view of the study of interaction mechanisms in a high intensity regime. Numerical codes were employed to simulate plasma hydrodynamics and atomic physics processes. A detailed description of the
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physics applied in such codes will be given along with a discussion of the limits of applicability to the experimental conditions under investigation.

On the other hand, when the density scalelength is much larger than the laser wavelength, the laser light can drive plasma instabilities which can scatter laser light reducing absorption, or can degrade the symmetry of the implosion. In fact the filamentation and/or the self-focusing instability can occur in this region and enhance laser intensity non-uniformities, consequently affecting the uniformity of the ablation pressure. In addition, other detrimental laser-driven plasma instabilities, including stimulated Brillouin scattering, stimulated Raman scattering and two plasmon decay, can be efficiently activated by a local increase of intensity subsequent to the onset of filamentation or self–focusing. As will be discussed in this thesis, these processes make the physics of laser interaction with large underdense plasmas of crucial concern, as shown by the wide range of laser beam smoothing techniques which are presently being developed world-wide, aimed to control, and eventually suppress these instabilities.

The experimental investigation presented in this thesis was concerned with the production and the characterization of underdense plasmas by laser irradiation of thin targets. Thin foil targets were used to generate long scalelength plasmas which were fully characterised in terms of electron density, using interferometric measurements, and electron temperature using time-resolved X-ray spectroscopy. Novel X-ray imaging techniques will be described which allowed the first time-resolved X-ray images of the plasma to be acquired, with a temporal resolution of \(\approx 100\) ps. Advanced numerical techniques were implemented and used in the analysis of the interferograms to obtain fully three dimensional density profiles of the plasma. An accurate comparison between the experimental electron density profiles and the predictions of the 1D hydrodynamic simulation will be presented and discussed. Measurements of the temporal evolution of the plasma electron temperature will be studied and the role of opacity in X-ray spectra will be analysed using numerical simulation. Moreover, the suitability of these preformed plasmas for interaction studies relevant to ICF will be discussed.

The interaction of a delayed laser pulse with these plasmas was studied under controlled conditions in a ICF-like regime. The effect of the delayed pulse on the plasma conditions was studied experimentally. In particular, the results of an experimental investigation on filamentation processes are presented and discussed in this thesis, which provide evidence of the onset of whole-beam self-focusing, even in presence of laser beam smoothing. The experimental data are modelled using the
kinetic theory of filamentation and the role of absorption processes and electron heat transport processes in the instability is taken into account.

A more direct study of the physics of absorption mechanisms and heat transport phenomena was carried out in a picosecond laser-plasma interaction regime where such processes can be more clearly identified. When a picosecond laser pulse is focused onto a solid target, a hot plasma is generated at electron densities well above the critical density. Consequently thermal conduction is expected to play a key role in the transport of energy from the laser absorption region to this high density plasma region. A better understanding of thermal transport mechanisms is therefore a key issue in the study of these plasmas.

The interaction of picosecond laser pulses with micron-sized plasmas was investigated in a high intensity regime. Short scalelength plasmas were produced by irradiating solid targets as well as layered targets with high power, prepulse-free picosecond laser light. Absorption measurements showing the effect of polarization and angle of incidence will be presented which indicate the predominance of resonance absorption mechanisms in the laser-target interaction. Detailed numerical modelling of both collisional and resonance absorption will also be discussed. Newly developed techniques for picosecond time-resolved X-ray spectroscopy will be described which enabled the first observation of transient heat propagation in solid matter in a picosecond regime. A strong departure of the measured mass ablation rate from the prediction of the classical heat transport theory is observed in the high intensity regime. Both absorption measurements and mass ablation rate measurements provide information on the interplay between absorption processes and electron heat transport. In agreement with two-dimensional Fokker-Planck hydrodynamic simulation, these measurements consistently indicate that strong lateral heat flow may occur in the subcritical region when the size of the laser focal spot becomes comparable to the longitudinal density scalelength.
A review of the main results of the fluid description of a plasma will be given in this Chapter, with emphasis on the propagation of electromagnetic waves in homogeneous as well as inhomogeneous plasmas. Electron plasma waves and ion-acoustic waves will also be considered, due to their relevance to the physics of laser-plasma interaction processes. The main damping processes for both electron and electromagnetic waves are taken into account. Finally, the propagation of light in steep density gradients is studied, in view of an experimental investigation of laser interaction processes with micron-sized plasmas presented in this thesis.
1.1 - Fluid Description of a Plasma

From a classical viewpoint, collisionless plasmas can be described (Chen, 1974) by the Vlasov equation which gives the temporal evolution of the phase space distribution function of each plasma species. Maxwell equations complete the description linking electric and magnetic fields with charge and current densities. Since such a detailed description is usually impracticable, a more tractable set of equations can be obtained averaging over the particle velocity and therefore treating the plasma as a fluid. In other words one focuses the attention on the macroscopic properties of the plasma, namely the thermodynamic state and the fluid velocity, which are described by velocity moments of the Vlasov equation. One basic condition for this description to be applicable is that the mean free path of the particles of the fluid is small compared to the characteristic lengths of the system.

Although a complete description of a plasma as a fluid would again require the solution of an infinite set of coupled moment equations, some approximations can be made which allow basic plasma features to be studied (Krueer, 1988). Most plasma properties can in fact be described by taking into account only the first two moments of the distribution function of each plasma species, i.e. the Continuity Equation and the Equation of Motion, provided that some assumptions can be made on the heat flow which lead to a suitable Equation Of State (EOS) for each plasma species. The EOS provides a relationship between the main state variables, i.e. pressure, density and temperature.

Two species plasma

For a collisionless plasma consisting of two species only, namely ions and electrons, the continuity equation and the equation of motion for the mean velocity \( u_j \) of particles with number density \( n_j \), charge \( q_j \), and mass \( m_j \), in the presence of an electric field \( E \) and a magnetic field \( B \), in the laboratory frame, are given by

\[
\frac{\partial n_j}{\partial t} + \nabla \cdot (n_j u_j) = 0 
\]

(1.1.1)

\[
m_j n_j \left( \frac{\partial u_j}{\partial t} + (u_j \cdot \nabla)u_j \right) = n_j q_j \left( E + \frac{u_j \times B}{c} \right) - \nabla p_j 
\]

(1.1.2)
where \( c \) is the speed of light and \( j = e, i \) for electrons and ions respectively. The electric and magnetic fields are related to the charge density, \( \rho_e = \sum n_j q_j \) and the current densities, \( \mathbf{J} = \sum n_j q_j \mathbf{u}_j \), by the Maxwell equations (Gaussian units)

\[
\nabla \times \mathbf{E} = -\frac{1}{c} \frac{\partial \mathbf{B}}{\partial t} \tag{1.1.3}
\]

\[
\nabla \times \mathbf{H} = \frac{1}{c} \frac{\partial \mathbf{D}}{\partial t} + \frac{4\pi}{c} \mathbf{J} \tag{1.1.4}
\]

\[
\nabla \cdot \mathbf{D} = 4\pi \rho_e \tag{1.1.5}
\]

\[
\nabla \cdot \mathbf{B} = 0 \tag{1.1.6}
\]

The displacement vector, \( \mathbf{D} \) and the magnetic induction vector \( \mathbf{B} \) are expressed in terms of the electric field \( \mathbf{E} \) and the magnetic intensity \( \mathbf{H} \) respectively, by the following constitutive relations

\[
\mathbf{D} = \varepsilon \mathbf{E} \tag{1.1.7}
\]

\[
\mathbf{B} = \mu \mathbf{H} \tag{1.1.8}
\]

\( \varepsilon \) and \( \mu \) being the dielectric permittivity and the magnetic permeability respectively.

**Plasma equation of state**

The equation of state plays a crucial role in the hydrodynamic motion of the plasma and it is also closely related to ionization processes. The particular choice of the EOS for an equilibrium plasma will give, due to thermodynamic consistency considerations, a constraint on the temperature-density dependence of the ionization state. The simplest EOS (More, 1991) can be derived using the assumption that electrons and ions give independent contributions to the plasma pressure and energy. In other words, the interaction between these two subsystems is neglected in this approximation. The total plasma pressure is therefore the sum of the pressure due to the electron gas and that due to the ion gas.

It can be shown that, for plasma temperatures above 10 eV, the electron contribution to the plasma pressure is dominant over the ion one. The plasma pressure is therefore determined by the electron gas pressure and, in the limit of the partially ionized ideal-plasma, one obtains the EOS given by Eq.1.1.9. It has been shown (More, 1991 and references therein) that this choice is consistent with an
1.1 Fluid Description of a Plasma

An ionization model which can be interpreted as a simplified version of the Saha equation (see Eq.2.2.2), where only ground-state ions are taken into account.

\[ p \cong p_e = n_e K_B T_e. \quad (1.1.9) \]

The applicability of Eq.1.1.9 to the particular experimental conditions under investigation will have to be evaluated on the basis of the expected electron density and temperature. In fact, at low temperatures and/or high densities, the interaction energy between the electrons becomes comparable to the thermal energy and the ideal-gas model breaks down. In addition the electrons become degenerate and the Maxwellian distribution function must be replaced by the Fermi-Dirac distribution. This circumstance will be considered in more details in Sect.3.1 where hydrodynamic modelling of high density plasmas is discussed.

For the purpose of modelling low density coronal plasmas for interaction studies (see Chpt.2 and Chpt.4), Eqs.1.1.1-9 provide a complete set of equations which can be used to describe the hydrodynamic evolution of the plasma as well as its ability to support propagation of electromagnetic waves and collective modes of oscillation. In particular circumstances a physical process occurring in a plasma, such as the propagation of e.m.waves or electrostatic waves, will give rise to changes in the plasma properties such as temperature, density and pressure which can be described using simple thermodynamic relations. If the time-scale of heat propagation is fast compared to the typical time-scale of the physical process under investigation, so that the temperature can be considered uniform over the fluid, the plasma is treated as isothermal. On the contrary, when the heat flow can be neglected the plasma is treated as adiabatic. In the intermediate regime no simple assumptions can be made and plasma properties cannot be derived simply by the first two moment equations as details of the distribution function, which are not taken into account by the first two moment equations, cannot be neglected.

In the study of laser interaction with plasmas one often deals with wave-like phenomena such as e.m., electron and ion waves. One can determine which transformation applies to the particular circumstance by comparing the propagation velocity of the wave with the electron thermal velocity. If \( \omega / k \ll v_{th} \), \( \omega \) and \( k \) being the frequency and the wave-number of the wave respectively, \( v_{th} = \sqrt{K_B T_e / m_e} \) being the electron thermal velocity, \( T_e, K_B, m_e \) being the electron temperature, the Boltzmann constant and the electron mass respectively, then the plasma undergoes an isothermal transformation and Eq.1.1.9 reduces to

\[ p_e = C n_e, \quad (1.1.10) \]
C = K_b T_e being a constant. On the contrary, when heat flow can be neglected as in the case of \( \omega/k >> v_{th} \), then an adiabatic transformation takes place and the relation between the pressure and the density becomes

\[
p_e = C n_e^\gamma,
\]

(1.1.11)

\( C \) being a constant and \( \gamma \) being the ratio of the specific heats, \( C_p/C_v \), given by \( \gamma = (2 + N)/N \), for particles with \( N \) degrees of freedom.

The equations given so far can be used to solve typical physical problems of laser-plasma interaction in the approximation of collisionless plasmas. A summary of the basic results of the theory of propagation of waves in plasmas will be given in Sect.1.2. In the case where collisions cannot be neglected, i.e. in low temperature and/or high density plasmas, Eq.1.1.2 must be modified to account for collisions. The effects of collisions can be included in this theory by adding a collision term to the Vlasov equation which leads to the Fokker-Planck equation. The effect of collisions on the propagation of waves in a plasma will also be considered in Sect.1.2.
1.2 - Wave Propagation in Homogeneous Plasmas

The set of equations given in the previous section can be used to study the propagation of charge/field perturbations in a plasma. In this section we will consider the so called two-fluid model which takes into account a two species plasma, i.e. electrons and one species of ions.

**Electron plasma waves**

Using an adiabatic equation of state of the type given by Eq.(1.1.11) and considering the ions at rest, it can be found that a plasma can support longitudinal electron plasma waves, also called plasmons, with wave-vector $k$ and angular frequency $\omega$ which, provided that $k \lambda_D << 1$, are linked by the dispersion relation

$$\omega^2 = \omega_p^2 + 3k^2 v_{th}^2, \quad (1.2.1)$$

$v_{th}$ being the electron thermal velocity, $k$ being the modulus of the wave-vector $k$ and $\omega_p$ being the characteristic plasma angular frequency

$$\omega_p = \sqrt{\frac{4\pi e^2 n_e}{m_e}} = 5.64 \times 10^4 n_e^{1/2} \text{ rad / sec} \quad (1.2.2)$$

where $e$ is the electron charge, $m_e$ is the electron mass and the electron density is expressed in cm$^{-3}$. A plasma with a given $n_e$ can therefore support longitudinal electron plasma waves of angular frequency $\omega$ greater than $\omega_p$ given by Eq.1.2.2.

**Ion acoustic waves**

Ion density perturbations can also propagate in a plasma with ion temperature $T_i$. Taking into account that ions are massive and that the response time is slow compared to that of the electrons, an isothermal EOS can be used which leads to the dispersion relation of the so called longitudinal ion-acoustic waves

$$\omega^2 = k^2 v_s^2, \quad (1.2.3)$$

where $v_s = \sqrt{(Z k_B T_e + 3k_B T_i) / m_i}$ is the sound velocity, $Z$ is the charge state and $m_i$ is the ion mass. These waves play a crucial role in the laser-plasma
interaction regime of interest for inertial confinement fusion as they give rise to one of the most important scattering process known as stimulated Brillouin scattering. This aspect will be discussed in detail in Sect.2.3.

**Electromagnetic waves**

Finally, looking for wave-like solutions of the linearized force equation (Eq.1.1.2) for the electric field in a plasma one obtains the dispersion relation for the propagation of transverse electromagnetic waves

$$\omega^2 = \omega_p^2 + k^2 c^2,$$  \hspace{1cm} (1.2.4)

where $c$ is the speed of light in vacuum. According to this equation, the propagation of an e.m. wave with angular frequency, in a plasma, can take place as long as $\omega > \omega_p$, the local characteristic plasma angular frequency. Consequently Eq.1.2.2 sets an upper limit to the electron density, the so called critical density or cut-off density, which can support the propagation of a transverse e.m. wave

$$n_{cr} = \frac{m_e \omega_p^2}{4\pi e^2} \cong 1.1 \times 10^{21} \lambda(\mu m)^2 \text{ cm}^{-3} \quad (1.2.5)$$

where $\lambda(\mu m)$ is the vacuum wavelength of the e.m. wave. It can easily be verified that such results are consistent with the following definitions of high frequency plasma conductivity $\sigma$, and plasma dielectric permittivity $\varepsilon$,

$$\sigma = \frac{i \omega_p^2}{4\pi \omega} \quad \varepsilon = 1 - \frac{\omega_p^2}{\omega^2}. \quad (1.2.6)$$

Eqs.1.2.1-6 summarize the results of the collisionless plasma theory most relevant to laser-plasma interaction. According to these results, a collisionless plasma with a given electron temperature and electron density can support three types of wave, that is longitudinal electron waves, longitudinal ion-acoustic waves and transverse e.m. waves.

**Collisionless damping of electron waves**

Since we will be also concerned with energy transfer mechanisms, it is crucial to know how the energy associated with these waves can be converted into thermal energy of the plasma. From a physical viewpoint it is clear that collisions between ions and electrons, which have been neglected so far in the collisionless theory, will
contribute to convert the energy of electrostatic and e.m.waves into thermal energy of the electrons. However, even in the absence of collisions, electrostatic waves can be damped via a collisionless process known as Landau damping. Collisionless damping of electron plasma waves can be regarded as a resonant mechanism which results in the acceleration of electrons whose velocity is close to the phase velocity of the electron wave. In fact, if this condition is satisfied, such electrons will experience the E-field associated with the wave as a constant accelerating field and will therefore gain energy from the wave. It can be found (Hughes, 1979) that strong Landau damping of an electron wave occurs when its phase velocity is

\[
\frac{\omega}{k} \leq 3 v_{th}, \tag{1.2.7}
\]

where \(\omega\) and \(k\) are the angular frequency and the wave-number of the electrostatic wave related to each other by Eq.1.2.1 and \(v_{th}\) is the electron thermal velocity.

The damping rate, \(1/\tau_L\), that is the rate at which the wave energy is transferred to the electrons gas, in the case of a Maxwellian electron distribution function, is given by

\[
\frac{1}{\tau_L} = \frac{\pi}{8} \frac{\omega_p^2 \omega^2}{k^3 v_{th}^3} \exp\left(-\frac{\omega^2}{2 k^2 v_{th}^2}\right). \tag{1.2.8}
\]

The Landau damping rate, as a function of the electron density, is plotted in Fig.1.2.1, for the case of two plasma frequencies relevant to laser-plasma interaction experiments with a Nd laser light.

![Fig.1.2.1. Landau damping rate of electron plasma waves at two frequencies relevant to laser-plasma interaction with Nd-laser light, i.e. the fundamental frequency and half of the fundamental frequency of the laser e.m.wave.](image)
As discussed in Sect.2.3, a number of non-linear laser-plasma interaction processes, including the two plasmon decay and the stimulated Raman scattering can give rise to production of electron plasma waves, typically at half of the laser angular frequency, that is \( \omega = 8.94 \times 10^{14} \) rad/s. In addition, in the presence of the critical density layer, linear conversion of laser energy into electron plasma waves, at the laser angular frequency, \( \omega = 1.79 \times 10^{15} \) rad/s can occur.

Furthermore, Landau damping of electron plasma waves leads to the production of energetic electrons, also referred to as hot electrons, which, typically, can have energies of up to several tens of keV. These electrons play a important role in the direct drive inertial confinement fusion scheme as they can propagate into the capsule, giving rise to pre-heating of the nuclear fuel, thus shifting the compression curve to an higher adiabat. More details on the production of hot electrons by electron plasma waves will be given in Sect.3.2.

**Collisional damping processes**

The effects of collisions on electrostatic and e.m.waves can be included in the theory introduced so far, by adding a collision term to the Vlasov equation thus obtaining the Fokker-Planck equation. From the point of view of moment equations, no change is expected in the continuity equation if ionization and recombination processes are neglected. On the other hand, in the case of a two component fluid, electron-electron collisions and ion-ion collisions lead to no net change of momentum. On the contrary, electron-ion collisions provide another efficient damping mechanism for electron plasma waves and can also account for damping of e.m.waves. In the case of electrostatic and e.m.waves the ions can be treated as fixed and therefore, the collision term in the equation of motion depends upon the electron velocity only.

From a physical viewpoint we can assume that the rate of change of momentum of the electrons is proportional to the momentum itself, the proportionality constant being the electron-ion collision frequency, \( \nu_{ei} \) relevant to this particular process. The total rate of change of momentum to be subtracted from the RHS of Eq.1.1.2 is therefore given by the rate of change of the momentum of the electrons

\[
\sum_{k \neq j} \left( \frac{\partial}{\partial t} n_j u_j \right) = \frac{\partial}{\partial t} n_e u_e = \nu_{ei} n_e u_e. \tag{1.2.9}
\]
By solving again the set of equations with this collision term included one finds a new dispersion relation for e.m. wave in a plasma

\[ \omega^2 = \omega_p^2 \left( 1 - \frac{\nu_{ei}}{\omega} \right) + k^2 c^2. \]  

(1.2.10)

This equation differs from Eq.1.2.4 by the imaginary term which accounts for dissipation. In the limit of large laser frequency, compared to the electron-ion collision frequency, this term vanishes and Eq.1.2.4 is recovered.

In the presence of an oscillating E-field, as in the case of an electrostatic wave or an e.m. wave, electrons are driven by the field and oscillate with a typical quiver velocity, \( v_q = eE/(m_e \omega) \). Electron-ion collisions can transform coherent oscillation energy of electrons into random (thermal) motion. The rate of energy dissipation due to collisions is simply given by the electron-ion collision frequency, \( \nu_{ei} \), times the energy density of electrons oscillating in the E-field of the wave, \( W_{os} = n m_e v_q^2/2 \). The rate of energy loss suffered by the wave is then \( W_E/\tau_{col} \), where \( 1/\tau_{col} \) is the collisional damping rate of the wave and \( W_E = E^2/8 \pi \) is the energy density associated with the E-field of the wave. Energy balance gives the collisional damping coefficient of the wave in terms of the angular frequency of the wave considered, electrostatic or e.m., and the electron-ion collision frequency

\[ \frac{1}{\tau_{col}} = \nu_{ei} \frac{\omega_p^2}{\omega^2}, \]  

(1.2.11)

where Eq.1.2.2 has been used to introduce the characteristic plasma frequency \( \omega_p \).

In the case of an e.m. wave with angular frequency \( \omega \) propagating in a plasma with electron density below the critical value, the damping process is often referred to as Inverse Bremsstrahlung (IB) as it can be considered the inverse process of bremsstrahlung emission due to scattering of the electrons by the ions.

The IB absorption coefficient is easily obtained as the damping rate given by Eq.1.2.11 divided by the propagation velocity of the wave energy. In the case of an e.m. wave the energy propagates at the group velocity given by \( v_g = d\omega/dk = c\sqrt{\epsilon} \), where \( \epsilon \) is the dielectric permittivity given by Eq.1.2.6. Using the classical expression for the electron-ion collision frequency

\[ \nu_{ei} = 2.91 \times 10^{-6} Z n_e \ln \Lambda \Theta_e^{-3/2} \text{sec}^{-1}, \]  

(1.2.12)
one finally obtains the relation which gives the absorption coefficient for an
em. wave with angular frequency $\omega$, travelling in a plasma with electron density $n_e$, 
electron temperature $\Theta_e$ and charge state $Z$,

$$\kappa_{ib} = 3.1 \times 10^{-7} Z n_e^2 \ln \Lambda \Theta_e^{-\gamma_2} \omega^{-2} \left(1 - \frac{\omega_p^2}{\omega^2}\right)^{-\gamma_2^2} \text{cm}^{-1}.$$ (1.2.13)

In this expression, $\ln \Lambda$ is the Coulomb logarithm, where $\Lambda = \frac{r_{\text{max}}}{r_{\text{min}}}$ is the ratio between the maximum and the minimum impact parameter relative to the electron-
ion scattering process. Under the conditions of interest, laser produced plasmas typically have $\ln \Lambda \approx 10$. However, in the calculations presented in this thesis, the results of a detailed study on this parameter (Skupsky, 1987) will be considered.

The importance of collisional absorption in laser-plasma interaction processes can be estimated by comparing the absorption length with the typical plasma scalelength. In typical ICF-like conditions, for example, the plasma is characterized by electron temperatures and densities of $\Theta_e \approx 1 \text{keV}$ and $n_e \approx 10^{20} \text{cm}^{-3}$ respectively. According to Eq.1.2.13, the absorption length for the fundamental harmonic of Nd laser light, in this case, is approximately $0.4 \text{ cm}$. Since this length is comparable to the typical density scalelength of the plasma, collisional absorption is expected to play a key role in the laser-plasma coupling in ICF-like conditions. Nevertheless, as will be discussed in Chapt.2, many other processes can also take place which can strongly affect collisional absorption. On the other hand, a more realistic approach to absorption processes in laser-produced plasmas requires that propagation mechanisms in inhomogeneous plasmas are taken into account. This aspect will be discussed in the next section.
1.3 - Inhomogeneous Plasmas

In order to provide a more satisfactory modelling of laser-plasma interaction, it is necessary to extend the theory of propagation of e.m.waves in plasmas to the more general case of inhomogeneous plasmas and, in particular, to plasmas whose electron density increases along the direction of propagation of the wave. This condition is relevant to most of the experimental conditions in which a laser beam interacts with a solid and generates a hot inhomogeneous plasma.

Weak density gradients and WKB approximation

We consider a plasma whose density, \( n(z) \) depends upon one spatial co-ordinate, \( z \) and an e.m.wave propagating along this direction. In the case of weak density gradients, i.e. when the density scalelength, \( L = n_e \left( \frac{\partial n_e}{\partial z} \right)^{-1} \) is much smaller than the local wavelength of laser light, \( \lambda(z) \), the solution can be found for any \( n(z) \) in the so called WKB (Wentzel, Kramers, Brillouin) approximation (Hughes, 1979). Looking for an oscillating solution of the wave equation for the electric field with a position dependent amplitude, \( \mathbf{E}(\mathbf{x}, t) = \hat{\mathbf{x}} E_x(z) \exp\left[ i(kz - \omega t) \right] \), where \( E_x(z) \), according to the WKB theory, is a slowly varying function of \( z \), and the electric vector has been chosen to be in the \( x \) direction, the amplitude of this electric field component depends upon the local plasma conditions according to the following expression

\[
|E_x(z)| = \frac{E_o}{\varepsilon(z)^{1/2}}, \quad (1.3.1)
\]

where \( \varepsilon(z) \) is the local value of the plasma dielectric function and \( E_o \) is the value of the incident E-field, i.e. the value of the electric field in the vacuum. According to Eq.1.3.1, the amplitude of the electric field increases as the light propagates towards higher densities. This effect can be simply explained using energy conservation considerations, as the energy flux at a given position \( z \), associated with the e.m.wave, must be conserved (neglecting absorption). By imposing this condition on the electric field one is led to compare the energy flux of the incident e.m.wave in the vacuum with the energy flux of the e.m.wave in the plasma.
The energy flowing through a given area perpendicular to the propagation direction must be conserved along the path, therefore one can write the following equation

\[ c \frac{\mathbf{E}_o \cdot \mathbf{E}_o}{4\pi} = \frac{d\omega}{dk} \frac{\mathbf{E}(z) \cdot \mathbf{E}(z)}{4\pi}, \]  

(1.3.2)

where \( c \) is the speed of light in the vacuum and \( d\omega/dk \) is the group velocity of the e.m. wave in the plasma. As the electron density increases, the group velocity decreases and the electric field consequently must increase according to Eq.1.3.1 in order to maintain the energy flux.

This swelling of the electric field in inhomogeneous plasmas can be particularly important in the interaction of laser light with steep density gradients, as in the case of picosecond laser pulses. In this case, the typical scalelength of the longitudinal density profile can be much smaller than the collisional absorption length, so that laser light can propagate up to the critical density layer without being severely damped. However, in this case, the condition for the validity of the WKB approximation is not strictly fulfilled since the density scalelength may be comparable to, or even shorter than the local laser wavelength. Furthermore, the WKB approximation breaks down in the proximity of the critical layer where both, the dielectric function \( \varepsilon \), and the wave-number of the e.m. wave \( k \), vanish.

**Steep density gradients**

A more accurate description can be obtained solving the wave equation for a particular choice of density profile. In the case of propagation of light in a linear density profile the solution of the equation for the electric field can be expressed analytically as a linear combination of the Airy functions. With the appropriate boundary conditions and assuming a density profile given by \( n_e = n_{cr} z/L \), one again finds that the electric field increases while the e.m. wave approaches the critical layer and reaches a maximum given by (Kruer, 1988)

\[ \left| \frac{E_x(z_{\text{max}})}{E_o} \right|^2 \approx 6.6 \left( \frac{L}{\lambda_o} \right)^{1/3}, \]  

(1.3.3)

where \( \lambda_o \) is the vacuum wavelength of the incident light. This maximum occurs at a distance from the critical density layer given by \( L - z_{\text{max}} \approx 0.3 \left( L \lambda_o^2 \right)^{1/3} \).

According to Eq.1.3.3, shorter wavelength light and longer density gradients will lead to a larger swelling of the electric field. However, as already observed above, the dependence upon the density scalelength should be considered in view of the
fact that collisional absorption will increase with $L$. Therefore maximum swelling will occur at an intermediate scalelength for which collisional absorption is not critical. For example, in the conditions relevant to the experimental investigation described in Chapt.5, the electric field of a 0.25 μm laser light, propagating in a 4 μm scalelength density gradient, will increase up to a factor of four with respect to the vacuum value. This maximum will take place at a distance of 0.2 μm from the critical density layer, where a substantial fraction of the incident light can propagate. It will be shown that this swelling effect can play a very important role in laser-plasma interaction and, in particular, in determining the interplay between different absorption mechanisms.

In order to evaluate the significance of such effects in laser-plasma interaction experiments, it is important to have a general knowledge of the plasma conditions which one can expect in a given experimental configuration. A preliminary investigation on the expected plasma conditions can be performed by using numerical simulation. A description of the techniques employed in this thesis to study hydrodynamic and atomic physics processes in plasmas will be given in the first two sections of the Chapt.2. This study will be mainly concerned with long-scalelength plasmas, while issues related to the investigation of short-scalelength, high density plasmas will be taken in account in Chapt.3.
CHAPTER 2

NANOSECOND LASER-MATTER INTERACTION

Long-scalelength plasmas are mainly considered in this Chapter. The hydrodynamics of plasmas produced by laser-matter interaction, in a nanosecond regime, will be investigated using one-dimensional numerical simulations. A summary of the main results concerning the atomic physics of such plasmas will be given with emphasis on equilibrium conditions and radiation emission properties. Finally, in view of the experimental study on the laser-plasma interaction physics presented in Chapt.4, an overview of the main results of the theory of laser-driven plasma instabilities will be presented, with emphasis on filamentation processes.
2.1 - Plasma Hydrodynamic Simulation

Basic hydrodynamic phenomena of laser-plasma interaction have been studied numerically using the one-dimensional Lagrangian code MEDUSA (Christiansen et al., 1974). A detailed description of this code has been given elsewhere (Rodgers et al., 1989). A summary of the physics included in the code will be presented in this section, with particular attention to those aspects of the simulation related to the hydrodynamics of long-scalelength plasmas. Sect.3.1 will focus on the application to short-scalelength high density plasma.

Lagrangian co-ordinates

The set of equations given in Sect.1.1 describes the fluid motion relative to the laboratory frame. However, the numerical implementation of this description for application to typical laser-matter interaction schemes is not straightforward. In fact, a number of difficulties arise from the presence of discontinuities when targets consisting of layers of different materials are considered. On the contrary, in the Lagrangian scheme, the reference frame is fixed within the fluid and the discontinuities are automatically preserved, making this method ideal to describe the hydrodynamic evolution of laser produced plasmas.

In the one-dimensional code MEDUSA, the plasma is described as a two component fluid, namely electrons and ions, characterized by four main variables, plasma mass density \( \rho(x, t) \), fluid velocity \( u(x, t) \), electron temperature \( T_e(x, t) \), and ion temperature \( T_i(x, t) \). Ions always behave as non-degenerate perfect gas, while electrons are described either by an ideal gas EOS or by a Thomas-Fermi EOS. The perfect gas may either be non-degenerate or partially degenerate or fully degenerate and a fully ionized plasma is assumed in this case.

For the particular circumstances of plasmas produced by nanosecond laser pulses, a perfect gas equation of state as given by Eq.1.1.9 will be typically assumed as it was found to satisfactorily model the plasma in the conditions achieved in this interaction regime. A more detailed analysis of degeneracy effects and a brief description of the Thomas-Fermi model for non-ideal electron gas will be given in Sect.3.1. Charge neutrality requires that the electron fluid and the ion
fluid share the same velocity. In this case the equation of motion given by Eq.1.1.2 assumes a much simpler form as the \textit{convective} derivative, $u \frac{du}{dx}$, necessary in the description in a fixed reference frame (Chen, 1974) is deleted. Neglecting the internal electric and magnetic fields, the equation of motion becomes

$$\rho \frac{du}{dt} = -\frac{dp}{dx},$$  \hspace{1cm} (2.1.1)

where $\rho$ is the mass density and $p = p_i + p_e$ is the hydrodynamic pressure. Finally the motion of the Lagrangian co-ordinate $r$ is determined by the fluid velocity $u$ according to the following equation

$$\frac{dr}{dt} = u(x,t).$$  \hspace{1cm} (2.1.2)

Each subsystem is also governed by an \textit{energy equation} which is obtained by balancing the rate at which energy enters the subsystem and the rate at which this energy goes into modifications of the thermodynamic and kinetic state according to

$$C_V \frac{dT}{dt} + B_T \frac{d\rho}{dt} + p \frac{dV}{dt} = S,$$  \hspace{1cm} (2.1.3)

where $S$ is the rate of energy input per unit mass, $C_V = (\partial U/\partial T)_{\rho}$ is the specific heat per unit volume, $B_T = (\partial U/\partial \rho)_{T}$ describes the variation of internal energy due to interaction between particles within the same subsystem, and $U = pV/(\gamma - 1)$ is the internal energy per unit mass. Electrons can exchange energy via thermal conduction, electron-ion collisions, bremsstrahlung emission and laser light absorption, while ions exchange energy via thermal conduction, electron-ion collisions, and viscous shock heating. Energy is exchanged between ions and electrons by means of electron-ion collisions at a rate given by Eq.1.2.12. Inverse bremsstrahlung absorption of laser light is modelled using the classical coefficients given by Eq.1.2.13. Thermal conductivity is modelled in terms of the classical Spitzer (Spitzer, 1953) conductivity with the thermal electron flux limited, for high temperature gradients, to the so called free-streaming limit, $F^{\text{max}}_e$ according to the following expression

$$\frac{1}{F_e} = \frac{1}{F_e} + \frac{1}{F^{\text{max}}_e},$$  \hspace{1cm} (2.1.4)

where the electron thermal flux, according to the usual definition, is proportional to the electron temperature gradient, i.e. $F_e = \kappa_e \nabla T_e$. 

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Hydrodynamic modelling of long scalelength plasma experiments

A target made of as many as three different layers can be specified, in the input conditions, in a planar, spherical or cylindrical geometry, by giving atomic number, atomic mass, mass density and thickness of each layer. Boundary conditions are specified by setting each target boundary to be fixed or free, according to the particular experimental configuration. Laser pulse parameters can be specified including wavelength, pulse-length, intensity, temporal shape and timing relative to the start-time of the hydrodynamic simulation.

Many physical processes can be included and controlled by means of other input logical switches. Resonance absorption can be included specifying the percentage of laser energy that, not absorbed by inverse bremsstrahlung and having reached the critical density layer, if any, is deposited at the critical density layer. Hot electrons can be included by setting the fraction of laser energy absorbed by the plasma via resonance absorption to be converted in hot electron energy. The hot electron temperature will be determined either as a function of the electron temperature at the critical density or as a function of laser intensity times the square of the laser wavelength. Ponderomotive force of the laser can also be included if required, in the momentum equation. Finally the ionization equilibrium can be calculated by the code using Saha ionization model.

Two laser-target configurations (see Section 4.1) will be mainly considered which will be referred to as Line Focus Geometry (LFG), and Cylindrical Geometry (CG) and a detailed description of these schemes is given in Sect.4.1. However, from the point of view of a 1-D hydrodynamic simulation of plasma production and expansion, there are only marginal differences between the two configurations. In fact, the two schemes mainly differ in the geometry of the delayed interaction pulse relative to the target. In the following, a brief summary of the main results of the hydrodynamic simulation for these two configurations will be given with particular attention to the spatial and temporal dependence of electron density and temperature.

In the CG (see Fig.4.1.2) the plasma was pre-formed by irradiating a 400 µm diameter, 500 nm thick Aluminium dot target, coated onto a 0.1 µm plastic substrate. Four 600 ps FWHM, 1.053 µm laser beams of the Vulcan laser at the Rutherford Laboratory (UK), were superimposed on target, two on each side, focused in a 600 µm diameter focal spot. With such a focal spot diameter the target was uniformly irradiated by the central region of the beams in order to limit perturbation effects due to possible cold plasma at the edges of the Al target.
In the LFG (see Fig.4.1.3) the target consisted of 700 µm long, 300 µm wide, 700 nm thick Aluminium layer coated onto a 0.1 µm plastic substrate. Four heating beams of the Vulcan laser were frequency doubled and superimposed on target, two on each side of the Al film. The focusing optics of each heating beam was tilted in order to generate a focal spot elongated along the longitudinal target axis, in order to optimize the matching between the laser focal spot and the target. In both cases the total heating irradiance was typically between $6.0 \times 10^{13}$ and $1.2 \times 10^{14}$ W/cm$^2$. As will be shown in Sect.2.3, this irradiance is below the typical threshold intensity (see Fig.2.3.1) for non-linear laser-plasma interaction mechanisms such as stimulated Brillouin scattering, stimulated Raman scattering and two plasmon decay. A description of these instabilities along with a summary of the main results relative to the dependence of the threshold intensities upon the main plasma parameters is presented in Sect.2.3.

In both configurations, the irradiation geometry was symmetric with respect to a plane parallel to the Al foil. Minor deviations from this symmetry can be expected due to the presence of the thin plastic substrate used to hold the Al foil. However, as shown below, in the conditions relevant to our experiment, such effects can be neglected. Therefore simulations were typically performed assuming a single-side irradiation, with a target thickness equal to half of the original thickness, and boundary of the target, opposite to the laser, kept fixed.

The simulation provides temporal evolution of the main hydrodynamic variables including electron and ion temperature, mass density, hydrodynamic velocity and pressure and average ionization relative to each cell in which the plasma is spatially sampled according to the numerical implementation of the Lagrangian scheme. Radiation losses due to bremsstrahlung emission are also accounted for although, for low and medium Z plasma, they were usually found to be negligible.

A portion of the energy not absorbed by inverse bremsstrahlung, typically 10-20%, was dumped at the critical density layer to simulate resonance absorption. This parameter was found to be not critical at all since, in the long scalelength regime under investigation here, most of the laser energy is absorbed due to collisional absorption, before the laser light reaches the critical density layer.

Although our attention will be mainly focused on plasma conditions at the time at which the delayed interaction beam reaches the preformed plasma, i.e. typically 2-3 ns after the peak of the heating pulse, the plasma formation and heating will also be considered in view of a comparison with the results obtained from experimental X-ray time resolved spectroscopy.
In order to evaluate the contribution of the thin plastic substrate to the hydrodynamic evolution of the plasma, numerical simulations of laser irradiated Al targets with and without the CH layer were carried out. A two-layer (CH-Al) target consisting of a 0.1 µm CH and a 0.25 µm Al was irradiated on the CH side, at an intensity of $3 \times 10^{13}$ W/cm$^2$. Fig. 2.1.1 shows a comparison between the electron temperature and density profiles obtained for this condition 2 ns after the peak of the pulse, and the analogous profiles obtained considering a 0.25 µm thick Al layer without plastic substrate.

According to the results of the simulation shown in Fig. 2.1.1, the plasma produced in the presence of the CH layer is expected to be slightly colder than the plasma produced in the case of irradiation of Al without CH. In fact, 2 ns after the peak of the heating pulse, the electron temperature is predicted to be 550 eV in the presence of the CH layer and 640 eV without the CH layer. This is essentially due to the fact that, early in the heating process, the interaction takes place in the plasma originating from the plastic layer, where collisional absorption is less efficient due to the lower charge state. As a consequence of the lower electron temperature, the plasma expands slightly less and the peak density, in the case of the CH coated target, is approximately 20% higher than the peak density of the plasma originating from the irradiation of the uncoated Al target.

In the actual experimental conditions, only one side of the Al foil is coated with plastic and consequently the plasma produced will be slightly asymmetric with
respect to the plane considered above. However, interaction experiments on this type of preformed plasmas (see Sec.4.6) were typically performed focusing the interaction beam on the side of the plasma originating from the bare Al. We will consequently focus our attention on this side of the plasma.

**Fig.2.1.2** shows the results of the simulation performed in the conditions of CG for an unsupported (without CH) Al target. The laser wavelength was 1.053 µm and the peak intensity was $3.0 \times 10^{13}$ W/cm$^2$. That was the typical value of the heating intensity, on each side of the Al foil, at which most of the experimental characterization was carried out. The electron temperature and density profiles relative to the peak of the heating laser pulses are shown in the graph in units of eV and cm$^{-3}$ respectively. The electron density was evaluated from the plasma mass density provided by the simulation according to the following relation

$$n_e = \rho Z^+/(A m_p), \quad (2.1.5)$$

where $Z^+$ is the ionization degree, $A$ is the atomic weight and $m_p$ is the proton mass, the difference between the mass of the proton and the mass of the neutron being neglected. Taking into account that the critical density at 1.053 µm is $n_{cr} \approx 10^{21}$ cm$^{-3}$, **Fig.2.1.2** shows that, at the peak of the pulse, the plasma is still overdense over a 50 µm long plasma column on each side of the target position.

**Fig.2.1.2.** Electron density and temperature profiles obtained from the hydrodynamic simulation at the peak of the heating laser pulse. A 250 nm thick Al target was irradiated at a wavelength of 1.053 µm and at an intensity of $3\times10^{13}$ W/cm$^2$. The electron conductivity was limited to 10% of the free streaming value.

The electron temperature is predicted to be uniform over the sub-critical region where it is expected to be as high as 1.2 keV, decreasing to approximately 0.8 keV...
into the higher density region, where laser light cannot propagate. Fig.2.1.3 shows the electron temperature and density profiles calculated using the same conditions as in Fig.2.1.2, but relative to 2.2 ns, 3.0 ns and 4.3 ns after the peak of the heating pulse. 2 ns after the peak of the heating pulse the plasma is fully underdense, the peak density being approximately $n_{cr}/10$ and the electron temperature, uniform over the whole plasma, being approximately 600 eV.

![Electron density and temperature profiles obtained from the same simulation of Fig.2.1.2 at three times after the peak of the heating laser pulse. A 250 nm thick Al target was irradiated at a wavelength of 1.053 µm and at an intensity of $3\times10^{13}$ W/cm$^2$.](image)

Hydrodynamic simulations were also performed for the experimental conditions of the LFG. In this case the laser wavelength was 0.527 µm and the target thickness used in the simulation was 350 nm while all the other parameters where left unchanged. According to Fig.2.1.4 and Fig.2.1.5, which give the electron density and temperature profiles at the peak of the pulse, and at the three times of interest considered above, the main difference between the two configurations under investigation is the slightly lower electron temperature and the higher electron density in the case of the LFG.

This is mainly due to the difference in the target thickness which results in more mass to be heated up by the laser. In fact, according to the simulation, more than 90% of the input laser energy is absorbed in both cases. At the peak of the laser pulse, the kinetic energy, related to the expansion velocity of the plasma, accounts for 54% of the absorbed energy, while the thermal motion of ions and electrons accounts for the remaining 36%.
Electron density and temperature profiles calculated at the peak of the heating pulse. A 350 nm thick Al target was irradiated at a wavelength of 0.527 µm and at an intensity of $3 \times 10^{13}$ W/cm$^2$. The electron conductivity was limited to 10% of the free streaming value.

Less than 2% of the input energy is lost by the plasma via bremsstrahlung emission. At 1 ns after the peak of the laser pulse, when the laser is off, the thermal energy is only 28% of the total energy while the kinetic energy accounts for 60% of the total energy, and energy losses by bremsstrahlung emission are still limited to approximately 2%.

Electron density and temperature profiles obtained from the same simulation of Fig.2.1.2 at three different times after the peak of the heating laser pulses. A 350 nm thick Al target was irradiated at a wavelength of 0.527 µm and at an intensity of $3 \times 10^{13}$ W/cm$^2$.

A comparison between the calculated electron density profiles shown above, and the experimental profiles, obtained from interferometric measurements, will be
Electron density profiles at 3.0 ns and 4.3 ns reported here will be directly compared with the experimental profiles taken at the same delay relative to the peak of the laser pulse (See Sect.4.3). This analysis will show the main limit of the simulation presented here, that is, the importance of two-dimensional expansion effects which are not accounted for by the 1-D hydro-code.

Furthermore, the temporal evolution of the X-ray emissivity on the plasma, obtained from time-resolved X-ray spectroscopy of Hydrogen and He-like Al emission, will be also compared with the calculation. In fact, as shown in Sect.4.5, the results of the hydrodynamic simulations presented in this section were post-processed in order to determine the X-ray emission properties of the plasma. A detailed analysis of atomic physics processes in plasmas is presented in the following section. In particular, the emission properties of plasmas are studied for the equilibrium conditions relevant to long-scalelength plasmas.
2.2 - Atomic Physics Processes in Plasmas

In the typical conditions relevant to laser produced plasmas, the peak of the spectral self-emissivity of the plasma is located in the X-ray region. A great deal of information on the physical properties of the plasma emitting region can be gained from a spectroscopic analysis of this emission. However, due to the complexity of the physical system under investigation, the analysis of experimental results is usually performed via comparison with numerical simulations.

A simplified steady-state model of the plasma is considered in the simulations, and the spectral features of the radiation emitted are studied as a function of a set of input plasma parameters including atomic number, ion and electron temperature, electron density and plasma size. These parameters can be set according to the experimental results or to the hydrodynamic simulation and will, in turn, allow for further information on the plasma to be gained and fed-back into the simulation.

The first issue to be addressed, in the study of emission of radiation from a plasma, is the kind of equilibrium to be considered. The degree of interaction among the three plasma sub-systems, namely electrons, ions and radiation, must be specified in order to determine the population of all available energy levels. Consequently, one can determine the spectral distribution of radiation energy emitted via bound-bound, bound-free and free-free transitions. Once the particular equilibrium is specified, the result of the numerical analysis of atomic physics processes can be compared with the experimental results provided that radiation transport effects and, possibly, time dependent effects are taken into account. In the following, a summary of the main results on plasma equilibria will be given with emphasis on the conditions to be fulfilled, in terms of density and temperature, for the various types of equilibria to hold. Some criteria which allow the validity of a steady-state approximation to be tested, will be given in Sect.3.4.

Thermal equilibrium (TE)

Although this kind of equilibrium does not apply to laboratory plasmas and is only approached in stellar interiors, it can be considered as a reference condition in the limit of high plasma density. A plasma is said to be in TE when electrons, ions and radiation are strongly coupled to each other and share the same temperature.
The populations $N_u$ and $N_l$ of two ionic bound levels, $u$ and $l$, with statistical weights $g_u$ and $g_l$ respectively, are given by the Boltzmann equation

$$\frac{N_u}{N_l} = \frac{g_u}{g_l} \exp\left(-\frac{\Delta E_{u,l}}{K_B T}\right), \quad (2.2.1)$$

where $\Delta E_{u,l}$ is the energy difference between the two levels and $T$ is the thermodynamic temperature of the plasma. The population of the ionization states is given by the Saha equation

$$\frac{N(Z+1)n_e}{N(Z)} = \frac{g_0(Z+1)}{g_0(Z)} \left[\frac{2\pi m_e K_B T}{\hbar^2}\right]^{3/2} \exp\left(-\frac{\chi(Z)}{K_B T}\right), \quad (2.2.2)$$

which gives the ratio between the population densities of two contiguous ionization states with charge $Z$ and $Z+1$, and statistical weights $g_0(Z)$ and $g_0(Z+1)$ respectively. The subscript “0” refers to the ground state of the ion and, in TE, it is, by far, the most populated one. In Eq.2.2.2 $\chi_0(Z)$ is the ionization potential of the ion with charge $Z$, $n_e$ is the electron density, $m_e$ is the electron mass and $\hbar$ is the Planck constant. The free electrons are distributed among the available energy levels and their velocity follows the Maxwell distribution function

$$dN_{v,v+dv} = 4\pi n_e \left(\frac{m_e}{2\pi K_B T}\right)^{3/2} \exp\left(-\frac{m_e v^2}{2 K_B T}\right)v^2 dv, \quad (2.2.3)$$

where $dN_{v,v+dv}$ is the number of electrons with velocities between $v$ and $v + dv$. Finally the spectral energy density of the radiation emitted by a plasma in TE is that of a black-body and is given by the Planck formula

$$u(v, T) = \frac{8\pi \hbar v^3}{c^3} \left(\frac{1}{\exp(h\nu/k_B T) - 1}\right), \quad (2.2.4)$$

Eqs.2.2.1-4 completely define the spectral properties of a plasma in thermal equilibrium. When plasma conditions are such that TE is not satisfied, a new set of equations will have to be derived.

The most important source of deviation from TE in a plasma is the relatively weak coupling of radiation with atoms and ions. If we assume that the populations of the available energy levels are entirely determined by particle collisions, radiation processes being ineffective, then we obtain another type of equilibrium called local thermal equilibrium (LTE). This equilibrium, in general, is still too restrictive for
laser produced plasmas to hold but, in some cases, can provide a simple estimate of the spectral emissivity, particularly in short-pulse laser produced plasmas where electron densities can be very high.

**Local thermal equilibrium (LTE)**

In an optically thin LTE condition, radiation can escape from the plasma and particle collisions by themselves account for the population of the ion energy levels. The main difference between LTE and TE is that radiation and particles do not share the same temperature. However Eqs.2.2.1-3 are still valid, provided that the thermodynamic temperature, $T$, is replaced with the electron temperature, $T_e$.

In contrast, the spectral properties of the radiation emitted by a plasma in LTE are now determined by bound-bound transitions which give rise to line radiation, and by free-bound and free-free transitions both yielding continuum radiation. In an optically thin plasma, a transition between an upper level, $u$, and a lower level, $l$ gives rise to a spectral line with an emitted power per unit volume given by

$$I_{u,l} = N_u A_{u,l} \Delta E_{u,l}, \quad (2.2.5)$$

where $A_{u,l}$ is the Einstein coefficient for the rate of spontaneous emission, $N_u$ is the upper level population density and $\Delta E_{u,l}$ is the energy difference between the two levels involved in the transition. Continuum radiation in an optically thin LTE plasma is due either to recombination (free-bound) or to electron-ion collisions (free-free or bremsstrahlung). The spectral power emitted, per unit volume, as continuum radiation, can therefore be written as

$$I(\nu) d\nu = n_e \sum_Z n_{i,Z} \left[ \gamma_{ff}(Z, T_e, \nu) + \sum_l \alpha_{fb}(Z, l, T_e, \nu) \right] h\nu \, d\nu, \quad (2.2.6)$$

where $\gamma_{ff}(Z, T_e, \nu)$ is the probability of an electron making a free-free transition in the proximity of an ion. This probability, in general, will depend upon the ion charge, $Z$, the electron temperature, $T_e$, and the frequency of the emitted photon. Similarly $\alpha_{fb}(Z, l, T_e, \nu)$ is the probability of a free electron making a transition to a bound level, $l$. The sum, in both terms, runs over all the ion states present in the plasma while, in the free-bound term, an additional sum runs over all possible final bound states.

As can be derived from the initial assumptions, LTE will be satisfied as long as collisional processes are dominant over radiative ones in determining the population
of the available energy levels. In other words, the probability of an excited ion decaying, due to spontaneous emission, must be much smaller than the probability of collisional decay. According to this condition, an operational criterion has been derived (McWhirter, 1965) based on the following inequality between collisional de-excitation and spontaneous emission

$$n_e n_u X(T_e, u, l) \geq 10 n_u A_{u,l},$$

(2.2.7)

where $X(T_e, u, l) = \{\sigma_{u,l} v_e\}$ is the electron de-excitation coefficient for a transition from an upper level, $u$, to a lower level, $l$, averaged over all possible electron velocities, $v_e$. Using a Maxwellian distribution, one finds the condition on the electron density for an optically thin plasma in order for LTE to hold

$$n_e \geq 1.7 \times 10^{14} \Theta_e^{1/2} \Delta E_{u,l}^3 \text{ cm}^{-3},$$

(2.2.8)

where the transition energy, $\Delta E_{u,l}$, and the electron temperature, $\Theta_e$ are expressed in eV. According to this condition, in a plasma with a given electron density and temperature, there is a maximum energy gap below which LTE holds. For example, in a plasma with an electron density of $10^{21}$ cm$^{-3}$ and an electron temperature of 500 eV, the maximum bound-bound transition energy for which LTE is satisfied is approximately 65 eV. In contrast, a bound-bound transition of 2 keV, typical of K-shell emission from H-like Al ions, can be considered in LTE if $n_e > 3 \times 10^{25}$ cm$^{-3}$, i.e. at densities well above the Al solid density.

It has been suggested (McWhirter, 1965) that if the plasma is optically thick, that is, the radiation is coupled with ions and electrons, this criterion may be relaxed to lower densities. In fact opacity effects give rise to radiative excitation which tends to balance radiative decay, making collisional processes effectively responsible for de-excitation. However, the condition given above should still be regarded as a necessary, but not sufficient condition for a plasma to be in LTE.

**Non-LTE plasmas**

In a plasma where the condition expressed by Eq.2.2.8 is not fulfilled, radiative processes are expected to play an important role in determining the population of the available energy states. In the general case all the possible electron processes will have to be taken into account and the populations of the levels will now depend upon atomic cross-sections. However, as these processes depend differently upon
2.2 Atomic Physics Processes in Plasmas

electron density, one can expect that, in particular conditions, some of them will be more efficient than others, leading to substantial simplification of the problem.

The electron processes of interest in this case are summarised in Table 2.2.I together with their symbolic representation and the dependence of the corresponding rates upon the electron density and the density of photons at the frequency involved in the transition.

<table>
<thead>
<tr>
<th>Process</th>
<th>Representation</th>
<th>D(n)</th>
<th>I(n)</th>
<th>D(\phi)</th>
<th>I(\phi)</th>
</tr>
</thead>
<tbody>
<tr>
<td>R Rad. Exc./De-exc.</td>
<td>(h\nu + N(Z) \Leftrightarrow N^*(Z))</td>
<td>-</td>
<td>-</td>
<td>(\rho h\nu)</td>
<td>-</td>
</tr>
<tr>
<td>R Free-Bound Trans.</td>
<td>(h\nu + N(Z) \Leftrightarrow N(Z+1) + e)</td>
<td>(n_e)</td>
<td>-</td>
<td>(\rho h\nu)</td>
<td>-</td>
</tr>
<tr>
<td>C Impact Exc./De-exc.</td>
<td>(e + N(Z) \Leftrightarrow e + N^*(Z))</td>
<td>(n_e)</td>
<td>(n_e)</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>C Impact Ion/3 b Rec.</td>
<td>(e + N(Z) \Leftrightarrow e + e + N(Z+1))</td>
<td>(n_e)</td>
<td>(n_e^2)</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>C Autoion./Diel. Rec.</td>
<td>(e + N(Z+1) \Leftrightarrow N^{**}(Z))</td>
<td>(n_e)</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
</tbody>
</table>

Table 2.2.I Symbolic representation of the main Radiative (R) and Collisional (C) electron processes taking place in a plasma which determine the population of all bound and free states available. The dependence of each rate upon electron density and photon density is also shown.

If we assume that the plasma is optically thin, then ion excitation and ionization will be supported by collisional processes. A question arises on how these processes are balanced, i.e., which of the possible de-excitation and recombination processes will have to be considered. In the general case, when density is not high enough to ensure LTE, all inverse processes will contribute to de-excitation and recombination. On the other hand, if the density is sufficiently low, all the collisional processes will become inefficient and only radiative de-excitation and recombination processes will have to be considered. In this last case the problem can be simplified leading to the so called coronal model.

Coronal equilibrium (CE)

In this case collisional excitation and ionization are balanced respectively by radiative de-excitation and recombination. By comparing the ionization and recombination rates one obtains the basic equation which governs ionization processes in the coronal equilibrium

\[ n_e N(Z) S(T_e, Z, o) = n_e N(Z + 1) \alpha_{fb}(T_e, Z + 1, o), \]  (2.2.9)

where \(S(T_e, Z, o)\) is the collisional ionization coefficient and \(\alpha_{fb}(T_e, Z + 1, o)\) is the radiative recombination coefficient. Eq.2.2.9 leads to the important result that, in
a coronal plasma, the population of the ionization states does not depend upon the
electron density and is given by

\[ \frac{N(Z+1)}{N(Z)} = \frac{S(T_e, Z, o)}{\alpha_{yb}(T_e, Z + 1, o)}. \] (2.2.10)

Once again it is assumed that most of the ions are in their ground state, \( o \), and
therefore both the ion populations and the coefficients in Eq.2.2.9 are relative to the
ground level. Also, it is assumed that electrons have a Maxwellian velocity
distribution as given by Eq.2.2.3 with a characteristic temperature, \( T_e \).

In a similar way one obtains the population \( N_u \) of an excited level, \( u \) by
balancing the collisional excitation from the ground level, \( o \), with the radiative de-
excitation onto a lower level, \( r \) and summing over all levels \( r < u \),

\[ n_e N(Z) X(T_e, g, u) = N_u(Z) \sum_{r<u} A_{u,r}. \] (2.2.11)

By substituting the population of the upper level given by Eq.2.2.11 into
Eq.2.2.5, we have an expression which gives the power emitted, per unit volume, in a
line relative to a transition between an upper level, \( u \), and a lower level, \( l \), in an
optically thin coronal plasma

\[ I_{u,l} = n_e N(Z) X(T_e, g, u) \Delta E_{u,l} \sum_{r<u} \frac{A_{u,l}}{A_{u,r}}, \] (2.2.12)

where \( X(T_e, g, u) \) is the collisional excitation coefficient and \( A_{u,l} \) is the Einstein
spontaneous emission coefficient. As already pointed out, line intensity in the
coronal equilibrium depends now upon the excitation and ionization coefficients,
and therefore upon atomic cross-sections relative to the particular transition.
Empirical expressions for these coefficients are available, which can be used to
obtain an analytical dependence of line intensities upon electron temperature.

The condition on the electron density, for coronal equilibrium to hold, can be
derived from the basic assumption that collisional de-excitation and recombination
processes must be inefficient when compared with the radiative processes. The
most probable collisional process in a plasma is that giving rise to a transition
between two neighbouring excited levels. This condition can be stated using the
criterion given by Eq.2.2.13 which ensures that the rate of spontaneous radiative
de-excitation of a level \( u \) is greater than the corresponding rate relative to the most
probable collisional de-excitation process leading to depopulation of the level \( u \).
It should be noted that the collisional de-excitation rate increases with increasing quantum number as an effect of the greater delocalization of the electron wave function which increases the probability of collisions.

\[ \sum_{r < u} A_{u,r} \geq n_e X(T_e, u, u - 1). \quad (2.2.13) \]

In contrast, the sum of the radiative spontaneous decay rates over all the lower available levels decreases, as an effect of the reduced oscillator strength of the most energetic transitions. Consequently, for a given density, there is always a quantum number, above which, the condition given by Eq.2.2.13 cannot be fulfilled.

Therefore one must define a further condition, from an operative point of view, according to experimental criteria. In fact, for diagnostic purposes of He-like and H-like plasmas, one is usually concerned with evaluation of spectral line intensities relative to transitions, on the ground state, from excited levels typically up to \( n_u = 6 \), that is, from the so called \( \alpha \) line to \( \varepsilon \) line. Assuming this level as the upper limit for which the condition given by Eq.2.2.13 is to be fulfilled, a quantitative evaluation of the criterion for hydrogen-like ions to be in coronal equilibrium is (McWhirter, 1965):

\[ n_e \leq 6 \times 10^{10} (Z + 1)^6 \Theta_e \frac{Y^2}{10 \Theta_e} \exp \left( \frac{(Z + 1)^2}{10 \Theta_e} \right) \text{cm}^{-3}, \quad (2.2.14) \]

where \( Z \) is the ion charge and \( \Theta_e \) is the electron temperature expressed in eV. In this case, the sum over the spontaneous emission coefficients was performed using tabulated values for hydrogen. The dependence of the inequality expressed by Eq.2.2.14, upon the quantum number of the excited level considered above, leads to the conclusion that, well above the upper limit \( n_u \), collisional excitation is balanced by collisional de-excitation, rather than by the spontaneous radiative decay, thus satisfying the condition for LTE to hold. The population of these higher states will therefore be given by Eq.2.2.1, where the thermodynamic temperature has, once again, to be replaced by the electron temperature, \( T_e \). In the middle region, i.e. for those levels close to \( n_u \), none of the equilibria discussed so far will be applicable. A new set of equations will have to be considered, as both collisional and radiative decay processes will contribute. Fig.2.2.1 shows the dependence of the maximum electron density for which the condition given by Eq.2.2.14 is satisfied, for hydrogen-like ions from He to Cu. The curve for the hydrogen-like He indicates that, for this plasma species, the electron density must be less than \( 10^{14} \) cm\(^{-3}\) for coronal equilibrium to hold. In contrast, due to the strong dependence of the
density limit upon the \( Z \) number of the species involved, hydrogen-like Al plasmas at electron densities up to approximately \( 5 \times 10^{18} \text{ cm}^{-3} \) can be considered in coronal equilibrium. This is a special case of particular interest in the study of long scalelength Al plasmas described in Chapt. 4. In fact, plasmas produced by laser irradiation of a thin Al target, after a few nanoseconds, reach the conditions of density and temperature given by Fig. 2.2.1 and can, therefore, be described using the coronal approximation (see Sect. 4.5).

**Fig. 2.2.1** Dependence of the maximum electron density for which the condition for coronal equilibrium is satisfied. Eq. 2.2.14 has been evaluated for hydrogen-like ions from Fluorine to Silicon and for Copper. The thicker solid curve has been calculated for hydrogen-like He and the corresponding density axis is given on the vertical axis in the RHS.

The curve relative to hydrogenic Cu should only be considered as indicative for the very high temperature limit. In fact, according to Eq. 2.2.10, the electron temperature required for hydrogen-like Cu to be the dominant ion species is (Mewe, 1987) approximately \( \Theta_e \approx 13 \text{ keV} \), whereas the dominant species at 1 keV is the Cu XXVI, that is Be-like Cu. In contrast, an Al plasma in coronal equilibrium becomes hydrogenic at approximately \( \Theta_e \approx 1 \text{ keV} \).

**Collisional-radiative equilibrium (CRE)**

If the electron density is neither low enough to satisfy the criterion (2.2.14) for coronal equilibrium, nor high enough for the criterion (2.2.8) for LTE, then we are in the case where both radiative and collisional decay must be taken into account (Van Der Sijde, 1990). On the hypothesis of optically thin plasmas, of all the processes listed in Table 2.2.I, only radiative excitation and ionization can again be neglected. The electrons still have a Maxwellian distribution as given by Eq. 2.2.3.
In the general case, a rate equation must be written for each bound level, taking into account all possible transitions, from and into this level. One would therefore obtain an infinite set of equations which would be impossible to solve unless some approximations could be made. The typical approach is to set, from case to case, an upper limit to the number of levels to be considered as the population of higher levels becomes very small and can eventually be neglected, without any substantial error being generated. On the other hand, as already observed, above some high quantum number level, collisional processes dominate over radiative ones, leading to LTE conditions. Once the truncation is operated, a finite set of equations is obtained which can be solved numerically using the available data-base on transition rate coefficients. This approach has in fact been considered in the analysis of most of the experimental X-ray spectra presented in Chapt.4 and 5. The atomic physics code RATION (Lee, 1984) was used to generate synthetic spectra and intensity ratios which were compared with experimental spectra.

The analysis presented so far in this Chapter provides the basic results which enable long-scalelength plasmas to be characterized as shown in Chapt.4. On the other hand, the interaction of laser light, using these long-scalelength plasmas, has also been studied in a regime relevant to ICF and the results are presented in Sect.4.6. The physics of laser interaction with long scalelength plasmas will be investigated in the following sections, where the most important laser induced instabilities are described. Particular emphasis will be given to filamentation processes, as they can seed the conditions for all the other instabilities to grow.
2.3 - Overview on Laser-Driven Plasma Instabilities

A summary of the basic results of the current theory on laser-driven plasma instabilities will be given in this section. In particular, the instabilities taken into account are those which are active in the laser interaction with long-scalelength underdense plasmas, as they have special relevance in inertial confinement fusion.

**Stimulated Brillouin scattering**

Stimulated Brillouin scattering can seriously affect the energy coupling between the laser and the plasma, as it can lead to a substantial scattering of the incident laser energy via excitation of ion plasma waves, according to the following conditions for matching of frequency, $\omega$, and wavevector, $k$,

$$\omega_o = \omega_{s,B} + \omega_i, \quad k_o = k_{s,B} + k_i,$$

(2.3.1)

where the indexes “$o$”, “$s,B$” and “$i$” refer to the incident and Brillouin-scattered light wave and to the ion-acoustic wave respectively. The minimum frequency of the light wave supported by the plasma for a given electron density $n_e$ is the local electron plasma frequency given by Eq.1.2.2. According to the dispersion relations for e.m.waves and ion-acoustic waves given in Sect.1.2, and considering that the frequency of the ion acoustic wave is very small compared to the laser frequency, one finds that the instability can occur throughout the underdense plasma.

From an experimental viewpoint we will be mainly concerned with the particular case of back-scattering, i.e. the scattering of light waves at an angle of 180 degrees to the incident laser wave-vector. The frequency of the back-scattered light is then given by (Giulietti et al., 1989 and references therein)

$$\omega_{s,B} = \omega_o \left[ 1 - \frac{2 v_s}{c} \left( \frac{1 - n_e/n_c}{1 + 4 k_o^2 \lambda_D^{-2} (n_e/n_c - 1)} \right)^{1/2} \right],$$

(2.3.2)

where $v_s$ is the speed of the ion acoustic wave defined in Sect.1.2 (see Eq.1.2.3), $\lambda_D = (K_B T_e / 4 \pi n_e e^2)^{1/2}$ is the Debye length, and $c$ is the speed of light. Of particular interest in the analysis of experimental results is the threshold intensity
2.3 Overview on Laser-Driven Plasma Instabilities

for the instability to occur. In the typical experimental conditions discussed in this work, laser light interacts with plasmas characterized by a longitudinal electron density gradient, \( L = n_e / (\partial n_e / \partial x) \), \( x \) being the direction of propagation of the laser light. In this case, the threshold intensity is determined by the density scalelength, rather than by the damping rates of the output waves as in the case of homogeneous plasmas. The SBS instability threshold in terms of the electron density scalelength \( L (\mu m) \), the plasma electron temperature \( \Theta_e (eV) \), and the laser light wavelength \( \lambda_o (\mu m) \), is (Baldis, 1991)

\[
I_{t,SBS} = 3.4 \times 10^{12} \frac{\Theta_e}{L \lambda_o} \text{ W/cm}^2.
\] (2.3.3)

Gradients in the expansion velocity of the plasma can also be effective in limiting the region of coupling (Krueer, 1988). Eq.2.3.3 can be modified in the presence of a gradient in the expansion velocity, replacing \( L \) with \( L_{V} = \varepsilon / (\partial \varepsilon / \partial x) \) is the velocity scalelength.

Stimulated Raman scattering

Another process which involves the scattering of light waves is stimulated Raman scattering. In this case an electron plasma wave with frequency \( \omega_{epw} \) and wave-vector \( k_{epw} \) is generated together with an e.m.wave with frequency \( \omega_{s,R} \) and wave-vector \( k_{s,R} \) according to the wave-matching conditions

\[
\omega_o = \omega_{s,R} + \omega_{epw}, \quad k_o = k_{s,R} + k_{epw}.
\] (2.3.4)

Following the same criterion considered above for the Brillouin scattering, the dispersion relations for the e.m.wave and electron plasma wave give a condition on the electron density for the instability to take place that is \( n_e \leq n_c/4 \). Again, in the case of inhomogeneous plasma and in the limit of \( n_e << n_c/4 \), the threshold intensity is (Baldis, 1991)

\[
I_{t,SRS} = 4 \times 10^{17} \frac{n_e}{L \lambda_o} \text{ W/cm}^2 \quad (n_e << n_c/4)
\] (2.3.5)

where \( L (\mu m) \) is the electron density scalelength and \( n_c \) is the critical density for the incident laser light of wavelength \( \lambda_o (\mu m) \). Eq.2.3.5 holds provided that the electron density is sufficiently small compared to the limit density at which the instability can occur, i.e. \( n_c/4 \), in order for the WKB approximation to hold (see
Sec.1.3). When this condition is not fulfilled a more accurate analysis must be used (Kruer, 1988) which predicts a reduction of the threshold intensity given by Eq.2.3.5 by a factor of $(\lambda_o/2\pi L)^{1/3}$.

Two plasmon decay

In this case the incident laser light decays into two electron plasma waves, the so called plasmons, with frequency $\omega_{epw,1}$ and $\omega_{epw,2}$ according to the following matching conditions

$$\omega_o = \omega_{epw,1} + \omega_{epw,2}, \quad \mathbf{k}_o = \mathbf{k}_1 + \mathbf{k}_2.$$ (2.3.6)

Considering that the frequencies of the two output plasma waves are approximately equal to the local electron plasma frequency, dispersion relations (Eqs.1.2.1-4) restrict the electron density region where the instability can occur to $n_e \cong n_c^4$. The inhomogeneous threshold intensity is given by (Baldis, 1991)

$$I_{t,TPD} = 5 \times 10^{12} \frac{\Theta_e}{L_{1/4} \lambda_o} \text{ W/cm}^2$$ (2.3.7)

where $\Theta_e ($eV$)$ is the electron temperature and $L_{1/4} ($\mu m$)$ is the electron density scalelength at $n_e \equiv n_c/4$, $n_c$ being the critical density for the incident laser light of wavelength $\lambda_o ($\mu m$)$.

A signature of the occurrence of two plasmon decay in laser plasma interaction experiments is the production of electromagnetic radiation at a frequency corresponding to half-integer harmonics of the incident laser light, and in particular at the three-half harmonics. This emission originates from the non-linear coupling of the incident e.m. wave with the plasma waves produced by the two plasmon decay instability. The spectral distribution of the three-half harmonic is strictly related to the dynamics of the coupling as well as the TPD instability itself. In particular, such a distribution can provide evidence (Giulietti et al., 1991) of propagation of the TPD plasmons in the density gradients, prior to their coupling with the incident laser light.

The threshold intensities for the three instabilities discussed here have been evaluated for the case of 1 µm laser light, according to the conditions given by Eq.2.3.3, Eq.2.3.5 and Eq.2.3.7. The result is shown in Fig.2.3.1 as a function of the density scalelength, in the range from 0.1 µm to 1 cm and for an electron temperature of 500 eV.
As already discussed above, SBS can lead to strong reflection of laser light and is therefore detrimental to ICF. TPD and SRS are also of particular concern to ICF as they both give rise to the production of electron plasma waves. In fact, as pointed out in Sect. 1.2, electron plasma waves can generate hot electrons via collisionless damping. These electrons are characterized by kinetic energies much higher than the electron thermal energy and, therefore, they can propagate into the inner part of the ICF fusion capsule where they can give rise to a pre-heating, making inertial confinement more difficult to achieve. The generation of hot electrons has in fact been related to the emission of three-half harmonics (Gizzi et al., 1992) by investigating the spectral distribution of X-ray bremsstrahlung emission at a photon energy $h\nu >> k_B T_e$.

As one can see from Fig. 2.3.1, the two plasmon decay and the Brillouin instability have a much lower threshold intensity compared to that of the Raman instability. In particular, for the density scalelength of interest for interaction studies related to inertial confinement fusion, typically $L \approx 1 \text{mm}$, the threshold for either instability is of the order of $10^{15} \text{W/cm}^2$. This intensity must therefore be regarded as an upper limit to the intensity of the laser light used to generate plasmas, in order to achieve an efficient energy coupling and to limit the generation of non-thermal electrons.

![Threshold Intensities](image)

**Fig. 2.3.1.** Threshold intensities as a function of the longitudinal scalelength of the electron density profile of the stimulated Brillouin scattering, stimulated Raman scattering and two plasmon decay instability. The threshold for the SRS for $n_e = 0.25 n_c$ is labelled with "*".

However this threshold must be further reduced in view of the fact that another important instability, known as the filamentation instability, can occur in the presence of laser and/or density non-uniformities. It has been shown that, even
when specifically designed techniques aimed to generate smoothed laser beams are used, this instability can still be active (Afshar-rad et al., 1992) and lead to local increase of the laser intensity well above the average intensity. A detailed analysis of this instability is given in the next section along with a discussion of the possible means by which the instability itself can be detected experimentally.
2.4 - Filamentation Instability

A summary of the current theoretical results concerning the filamentation instability (FI) will be given in this section followed by a detailed description of second harmonic and X-ray emission processes from laser-plasma filaments. Sect.4.6 will be devoted to the analysis of experimental results obtained using both second harmonic emission and X-ray emission as diagnostic tools to investigate the temporal and spatial features of filamentation processes.

Growth rate and heat transport phenomena

Filamentation of laser light in a plasma can occur when a small perturbation in the intensity profile of the incident laser beam induces a perturbation in the electron density. This perturbation can be generated either directly, via the ponderomotive force, or indirectly as a consequence of localized collisional heating and subsequent plasma expansion. Once the perturbation has been generated, refraction of the laser light in the electron density perturbation enhances the intensity perturbation providing the feedback for the instability.

The most important theoretical results concerning the filamentation instability have been derived using a simple model in which a sinusoidal intensity perturbation is imposed on a plane light wave interacting with a plasma. Recently, the effect of non-local heat transport has also been included (Epperlein, 1991) in the theory of laser filamentation in plasmas to account for those conditions where the classical Spitzer-Härm conductivity cannot be used to describe thermal electron transport, that is, when the electron temperature scalelength is shorter than the electron mean free path. In these conditions thermal transport phenomena must be described solving the electron Fokker-Planck equation. An effective conductivity $\kappa_{\text{FP}}$, was determined by studying the effect of a sinusoidal spatial modulation of wavenumber $k_\perp$ in the laser field, incident on an initially homogeneous plasma. It was found that, in the limit of large perturbation wavelengths, the effective conductivity $\kappa_{\text{FP}}$ converges to the Spitzer-Härm classical conductivity given by (Rose, 1993)

$$\kappa_{\text{SH}} = (14 n_e K_B T_e / m_e v_e)(1 + 3.3|Z|),$$

where this expression is consistent with the heat flow given by $q_{\text{SH}} = -\kappa_{\text{SH}} \nabla (K_B T_e)$. In contrast, in the small wavelength limit, flux inhibition occurs resulting in $\kappa_{\text{FP}} < \kappa_{\text{SH}}$. This is a consequence of non-
local transport driven by electrons with energies of typically several times the background thermal electron energy. Due to their higher kinetic energy and consequently lower collisionality compared to that of the thermal electrons, these hot electrons can propagate for several perturbation wavelengths, reducing the effectiveness of the thermal heat flow.

This behavior can be summarised considering the ratio of the effective conductivity to the Spitzer-Härm conductivity as a function of $k_{\perp}\lambda_s$, where $\lambda_s = (K_B T_e)^2 (4 \pi n_e e^4 \ln \Lambda)^{-1}(Z + 1)^{-1/2}$ is the electron stopping length which, from a physical viewpoint, is approximately (Berger et al., 1993) the geometric mean of the mean-free path for electron-electron scattering and the mean-free path for electron-ion scattering. The best fit to the data obtained from the numerical solution was found to be (Epperlein, 1990)

$$\frac{\kappa_{FP}}{\kappa_{SH}} = \frac{1}{1 + (30 k_{\perp} \lambda_s)^{4/3}}.$$  \hfill (2.4.1)

A plot of the ratio of Eq.2.4.1 as a function of $k_{\perp}\lambda_s$ is presented in Fig.2.4.1. According to this curve, non-local effects in the electron conductivity become important when the perturbation wavelength is smaller than $\approx 100\lambda_s$, where the effective electron conductivity is approximately 50% of the Spitzer-Härm value.

![Fig.2.4.1. Ratio of the effective electron conductivity to the classical Spitzer-Härm conductivity as a function of $k_{\perp}\lambda_s$, where $k_{\perp}$ is the wavenumber of a sinusoidal perturbation of the laser intensity incident on an initially uniform plasma and $\lambda_s$ is the electron stopping length.](image-url)
This modification to the electron heat conductivity leads to the following expression for the spatial growth rate of the filamentation instability

\[ K = \frac{k_{\perp}}{2\sqrt{\varepsilon}} \left[ \frac{2n_e}{n_c} \left( \gamma_p + \gamma_T \frac{k_{SH}^2}{k_{FP} k_{\perp}^2} \right) - \frac{k_{\perp}^2}{k_o^2} \right]^{1/2}, \quad (2.4.2) \]

where \( \varepsilon = 1 - n_e/n_c \) is the plasma dielectric function (see Eq.1.2.6) and \( n_e \) and \( n_c \) are the electron density and the critical electron density respectively. In the first term of Eq.2.4.2, \( \gamma_p = (1/4)(Z/(Z + 1))(v_q^2/v_{th}^2) \) accounts for ponderomotive effects, \( v_q \) and \( v_{th} \) being the quiver and thermal velocity of the electrons respectively, and \( Z \) is the charge state of the plasma. In the second term \( \gamma_T = c^2 S/\omega_0^2 k_{SH} K_B T_e \) accounts for thermal effects, \( S = W_e/\tau_{col} \) being the background inverse bremsstrahlung heating rate, with the collisional damping rate of the e.m.wave given by Eq.1.2.11, \( K_B \) being the Boltzmann constant and \( T_e \) the electron temperature. The third term is due to diffraction and gives negative feedback as it tends to defocus the beam. The important consequence of non-local electron transport in the theory of filamentation instability is that the threshold of the instability is substantially reduced since, generally speaking, the growth rate is increased. In addition, an optimum perturbation wavelength is found which maximises the growth rate, in contrast with the theory based on the classical electron transport, which predicts a constant growth rate over a wide range of perturbation wavelengths (Biancalana et al., 1993 and references therein).

More recently (Berger, 1993) it has been pointed out that when the perturbation wavelength is much smaller than the electron stopping length, i.e. \( k_{\perp}\lambda_s \gg 1 \), even the modified electron transport theory is not adequate and collisionless Landau damping, must be taken into account. The effective non-local conductivity given by Eq.2.4.1 must be replaced by the Landau damping conductivity

\[ \kappa_{LD} = \sqrt{\frac{2}{\pi}} \frac{v_{th} n_e}{k_{\perp}} \quad (2.4.3) \]

when \( k_{\perp}\lambda_{90} = k_{\perp}\lambda_s \sqrt{Z} \approx 1 \), i.e. when a thermal electron travels a distance comparable to the perturbation wavelength without undergoing a large angle (90 degrees) scattering collision. Although both conductivities have a very similar dependence on \( k_{\perp} \), for \( k_{\perp}\lambda_s \gg 1 \), one can find that the Landau conductivity increases more rapidly with \( Z \) than the non-local one by a factor \( \sqrt{Z} \). In other words for high \( Z \) plasmas Eq.2.4.3 predicts that the strength of thermal filamentation relative to short perturbation wavelengths is reduced with respect to the nonlocal
2.4 Filamentation Instability

electron transport case as a consequence of the more efficient dissipation arising from collisionless Landau damping.

Laser-plasma filaments and X-ray emission

The occurrence of filamentation and/or self-focusing of a laser beam in a plasma causes substantial modifications of the local plasma conditions, producing local density depression and an increase in the electron temperature. These circumstances, and the consequent effects on emission, absorption and scattering of radiation, can provide indirect evidence of the process itself.

An increase in the local electron temperature will affect the emission properties of the plasma. In particular the X-ray radiation emitted from the plasma region involved in the filamentation processes will be modified in terms of spectral distribution and intensity. All the emission processes discussed in Sect.2.2, arising from bound-bound, free-bound and free-free transitions, will reflect the modifications occurring in the plasma, with a time response typically of the order of a picosecond (see Sect.3.5). On the other hand, the time-scale of filamentation processes is of the order of the perturbation scalelength divided by the speed of sound (see Eq.1.2.3). Assuming a scalelength $\lambda_\perp \approx 10^{-3}$ cm and a speed of sound $v_s = 10^7$ cm/s, one finds a time-scale of $\approx 100$ ps. Since this time is much larger than the typical X-ray response time, X-ray emission can provide valuable information on the temporal evolution of the filamentation processes (Giulietti et al., 1991; Afshar-rad et al., 1992). Moreover, if the spatial distribution of this emission can be monitored, the heat transport phenomena can also be investigated. An experimental study based on these techniques will be presented in Sect.4.6.

Second harmonic emission from laser-plasma filaments

The onset of filamentation during the propagation of laser light in underdense plasmas has been related to emission of radiation at the second harmonic (SH) of the laser frequency. SH emitted forward, i.e. in the direction of the laser beam, (Biancalana et al., 1993) as well as at 90 degrees to the axis of the incident laser beam has in fact been studied in terms of both spatial (Stamper et al., 1985; Young et al., 1989; ) and spectral distribution (Giulietti et al., 1989, Giulietti et al.,1990).

In fact, in the presence of gradients perpendicular to the laser beam axis in the electron density distribution and in the laser intensity, a current oscillating at the second harmonic of the laser light is generated in the interaction region whose density, $J_{2\omega}$ is given by (Erokhin, et al., 1974; Deha et al.,1992)
2.4 Filamentation Instability

\[ J_{2\omega} = \frac{i e^3}{m_e^2 \omega^2} \left[ \frac{n_e^2}{4} \nabla (E_{\omega} \cdot E_{\omega}) + \frac{(\nabla n_e \cdot E_{\omega}) E_{\omega}}{\epsilon} \right], \quad (2.4.4) \]

where \( e, m_e \) and \( n_e \) are the electron charge, mass and density respectively, \( E_{\omega} \) is the electric field at the laser frequency \( \omega \), and \( \epsilon \) is the plasma dielectric function given by Eq.1.2.6. If we assume that gradients are related to filaments of length \( L \) and radius \( r \), electromagnetic radiation at the second harmonic is generated according to Eq.2.4.4, whose Poynting vector, at a distance from the filament large compared to \( L \), can be written in the following form

\[ S = A \frac{L^2}{R^2 \omega^4} E_{\omega}^4 n_e^2 r^2 \frac{R}{R} \frac{\sin^2(\chi)}{\chi^2} f(\psi), \quad (2.4.5) \]

where \( A \) is a function of universal constants only, \( R \) is the vector from the filament to the viewer forming an angle \( \psi \) with the axis of the filament and \( \chi = (L/2)(2 k_\omega - k_{2\omega} \cos \psi) \). The function \( f(\psi) \) has several maxima for \( \psi < 30^\circ \), but vanishes for \( \psi = 0 \). More details will be given in Sect.4.6 where experimental results concerning the interplay between second harmonic emission and filamentation will be presented.

It should be noted that, in particular circumstances, i.e. in the presence of reflected or back-scattered laser light, second harmonic emission perpendicular to the axis of the filament can be generated. This is very important when laser light interacts with underdense plasmas as, in this case, no other mechanisms can give rise to second harmonic emission at 90 degrees. In this case Eq.2.4.4 must be replaced by the following expression (Shen, 1984)

\[ J_{2\omega} = \frac{e^3}{2 i m_e^2 \omega} (E_{\omega} E_{\omega}^* + E_{\omega}^* E_{\omega, B}) \cdot \nabla n_e, \quad (2.4.6) \]

where \( E_{\omega} \) and \( E_{\omega, B} \) are the amplitudes of the electric field of the incident and back-scattered laser light and \( \nabla n_e \) is the density gradient in the plasma. According to this expression, radial density gradients are needed to generate second harmonic emission at 90 degrees. And this is usual for gradients generated by the filamentation instability. The back-scattered component of light at approximately the same frequency of the laser light required Eq.2.4.6 can be provided by stimulated Brillouin back-scattering. In this case the frequency of the second harmonic light is given by \( \omega_{SH} = \omega_0 + \omega_B \), where \( \omega_B \) is given by Eq.2.3.2. This process has indeed been observed experimentally (Giulietti et al., 1989) in the interaction of laser light with underdense plasmas by monitoring the temporal evolution of the spectral distribution of second harmonic radiation scattered by the plasma at 90 degrees.
CHAPTER 3

PICOSECOND LASER-MATTER INTERACTION

The physics of laser-matter interaction in a picosecond regime will be investigated in this Chapter. The hydrodynamics of high density plasmas produced by high intensity, picosecond laser pulses, interacting with solid targets will be examined and the role of non-ideal plasma effects will be studied. Laser-plasma coupling phenomena will be analysed and the role of polarization and angle of incidence will be discussed. Collisional absorption processes in a high intensity regime will also be taken into account. Finally, the validity of the steady-state atomic physics modelling presented in Chapt.2 will be discussed.
3.1 - High Density Plasma Effects

The hydrodynamic simulation of laser produced plasmas in the nanosecond regime was performed assuming an ideal gas equation of state for both ions and electrons. The ideal gas approximation is expected to be valid in this case as the electron densities of interest, typically of the order of $10^{21} \text{cm}^{-3}$, ensure that the electron-electron, as well as the electron-ion interaction energy is small compared to the electron kinetic energy. Therefore, the electron gas can be considered ideal. In this case, the equation of state for the non-degenerate electron gas is simply given by Eq.1.1.9. On the other hand, at sufficiently low electron temperatures or high electron densities, non-ideal effects start playing a role in determining the thermodynamic state of the plasma.

The most important non-ideal processes will be examined in this section, including electron degeneracy effects, electron-ion and ion-ion interactions. The conditions on plasma density and temperature will be derived for which such effects must be included in the hydrodynamic modelling of the plasma, as they give important contributions to the equation of state.

Quantum mechanical effects

This limit is reached when the kinetic electron energy, $K_B T_e$, becomes comparable to, or smaller than the Fermi energy, $K_B T_F$, where $T_F$ is the Fermi temperature given by

$$ T_F = \frac{\hbar^2}{2 m K_B} \left( \frac{3 \pi^2 n_e}{2} \right)^{2/3} \approx 3.64 \times 10^{-15} n_e^{2/3} \text{eV}, \quad (3.1.1) $$

where $n_e$ is the electron density in cm$^{-3}$. In these circumstances the Maxwellian electron distribution function for the electron velocities must be replaced by the Fermi-Dirac (FD) distribution function and the EOS will be consequently modified. In the case of a partially degenerate electron gas, that is when $T_e \approx T_F$, Eq.1.1.9 will be replaced by an appropriate expansion (Christiansen, 1974) of the Fermi-Dirac pressure, in terms of the degeneracy parameter $\zeta = T_e/T_F$. For a fully degenerate
High Density Plasma Effects

Electron gas the pressure becomes a function of the electron density only and is given by the Fermi pressure

\[
p_F = \frac{\hbar^2}{5m_e} \left( \frac{3\pi^2}{2} \right)^{2/3} n_e^{5/3} \approx 2.33 \times 10^{-33} n_e^{5/3} \text{ Mbar} \quad (3.1.2)
\]

where \(n_e\) is the electron density in cm\(^{-3}\). According to Eq.3.1.1, a plasma with an electron density of \(10^{21}\) cm\(^{-3}\) is fully degenerate for electron temperatures below 1.5 eV. Since, in the typical plasma conditions relevant to long scalelength plasma experiments, the electron temperatures of interest are typically greater than 100 eV, we can conclude that degeneracy effects of the electron gas can be neglected in this case, and Eq.1.1.9 is expected to be accurate.

In contrast, in the case of plasmas produced by picosecond laser pulses, conditions can be substantially different, as discussed in Chapt.5, and a more accurate evaluation is required. A simple estimate of the importance of degeneracy effects in such plasmas can be obtained by considering the maximum electron density which can be achieved in this interaction regime. As will be discussed in Sect.5.3, plasmas with electron densities in excess of \(10^{23}\) cm\(^{-3}\) and electron temperatures of several hundred eV were generated (Riley et al., 1992) in picosecond laser interaction with Al targets. On the other hand, higher densities are expected to be achieved, although characterized by lower electron temperatures. An upper limit to the electron density can reasonably be set at the solid density of fully ionized Al, that is \(7.8 \times 10^{23}\) cm\(^{-3}\). According to Eq.3.1.1, the corresponding Fermi temperature is \(T_F \approx 31\) eV which suggests that, although degeneracy effects are not dominant over the whole plasma, some contributions to the electron EOS can be expected in the cold supercritical density region. Nevertheless, as will be shown below, these contributions are not expected to significantly affect the hydrodynamic evolution of the plasma.

Coulomb interactions

There are however other sources of deviation from the ideal gas behaviour of ions and electrons which have to be accounted for at high plasma densities. In fact as the plasma density increases, the separation between the particles decreases proportionally to \(\rho^{-1/3}\) and, in the case of charged particles, a limit is reached when the Coulomb interaction energy, proportional to the inverse of the particle separation, becomes comparable to the kinetic energy, \(K_B T\).
In a two component plasma, consisting of ions and electrons, one can consider three typical inter-particle distances, namely the ion-ion, the ion-electron and the electron-electron distance. Each distance can be associated to a particular effect that may contribute the EOS of the plasma. In the case of the ions, for example, when their separation becomes smaller than the Debye screening length, given by
\[ \lambda_D = 743 \Theta_e^{1/2} n_e^{-1/2}, \]
Coulomb interactions between ions become important and the plasma is said to be strongly coupled.

In general, the importance of the inter-particle (electrons and/or ions) coupling due to the Coulomb interaction can be assessed by evaluating the coupling parameter, that is, the ratio between the Coulomb energy and the kinetic energy
\[ \Gamma = \frac{Z_1^* Z_2^* e^2}{R K_B T}, \] (3.1.3)
where \( Z_j^* \) is the charge of the \( j \)-th particle (ion or electron), \( R \) is the inter-particle separation and \( K_B T \) is the particle kinetic energy. This coupling parameter has been evaluated for a range of relevant plasma temperatures and densities. Fig.3.1.1 shows the contour curves in the temperature-density space relative to \( \Gamma = 1 \) for the three classes of interaction considered above, for a fully ionized Al plasma. The portion of the \((n, T)\) space below each curve corresponds to coupling parameters greater than unity, that is, the corresponding Coulomb energy is greater than the kinetic energy. The condition \( \zeta = 1 \) for a partially degenerate electron gas is also shown (dashed-dotted curve) for comparison. In this case, the portion of the \((n, T)\) space below the curve corresponds to the condition of a partially degenerate electron gas.

---

**Fig.3.1.1** Contour curves in the temperature-density space, relative to \( \Gamma = 1 \) for the three classes of interaction (e-e, e-i, i-i) for a fully ionized Al plasma. The portion of the \((n, T)\) space above each curve corresponds to coupling parameters greater than unity. The (dashed-dotted curve) represents the condition \( \zeta = 1 \) for partially degenerate electron gas.
The difference in the slope of the last curve, compared to the slopes of the curves relative to the Coulomb interaction, is due to the dependence of the Fermi energy upon the plasma density given by Eq.3.1.1, where \( E_F \propto n^{2/3} \), compared to the Coulomb energy, where \( E_C \propto n^{1/3} \). It should be noted that the calculation has been performed assuming a fully ionized plasma over the whole \((n, T)\) region shown. This will be strictly valid at high temperatures only. In the low temperature region, typically below 100 eV, this assumption is not fulfilled; the charge state of the ions decreases, leading to a weaker Coulomb interaction. The curves shown here in the low temperature region, are therefore to be considered an upper bound to the \((n, T)\) region where strong coupling occurs.

According to Fig.3.1.1, strong ion-ion coupling occurs in the typical plasma conditions of laser produced plasmas. However, as already discussed in Sect.2.1, the contribution of the ions to the plasma pressure and energy is small compared to that of the electrons and, therefore, even in the presence of strong ion-ion coupling, the effects on the EOS can be neglected. In contrast, the Coulomb coupling between ions and electrons (dotted curve) has important consequences on the EOS and therefore, cannot be ignored. In this case, a new EOS is derived by taking into account Coulomb interactions and ionization processes in a self-consistent treatment. Several models have been developed (Baker & Johnson, 1991) to deal with ionization processes in dense matter which are based on the Thomas-Fermi model (Feynman et al., 1949). A brief summary of the basic concepts behind this model are given below.

### Thomas-Fermi model

The Thomas-Fermi model is applicable when \( Z >> 1 \) and approaches the problem considering all the electrons (bound and free) as partially degenerate, with an energy given by the sum of the kinetic energy and the Coulomb potential energy

\[
E_e = \frac{p_e^2}{2m_e} - e \phi(r)
\]

where the potential \( \phi(r) \) varies continuously through the ion. The electron density \( n_e(r) \) also varies continuously through the ion subject to appropriate boundary conditions. In the proximity of the nucleus \( \phi(r) \rightarrow Ze/r \) and, at a distance from the nucleus equal to the ion sphere radius, \( \phi(R_o) = 0 \). The density and the potential are then obtained consistently by solving Poisson’s equation, using an iterative procedure and assuming a spherically symmetric electron density
3.1 High Density Plasma Effects

distribution around the atom. Finally, once convergence has been achieved, the electron pressure is taken as the Fermi degenerate pressure at \( r = R_0 \). In the limit \( K_B T_e \ll E_F \), when the Fermi energy is much greater than the electron kinetic energy, this model gives the same results as the Fermi degenerate electron gas discussed at the beginning of this section. In the opposite limit, when the electron kinetic energy is much greater than both the Fermi and the Coulomb energy, i.e. \( K_B T_e \gg E_F, E_C \), the ideal gas results are recovered.

Non-ideal plasma effects in the hydrodynamic simulation.

Hydrodynamic simulations of plasma production and evolution, in the picosecond regime, have been performed including non-ideal plasma effects in the calculation. The results of the simulation, performed assuming an ideal gas EOS for the electrons, are compared with the results obtained assuming a Fermi-Dirac EOS and a Thomas-Fermi EOS. A 12 ps, 0.25 \( \mu \)m laser pulse interacted with a solid Al target at irradiances in the range between \( 10^{14} \) and \( 10^{17} \) W/cm\(^2\). In fact, a detailed experimental investigation of interaction mechanisms, including absorption of laser light, heat transport phenomena and X-ray emission, was performed using these parameters. Fig.3.1.2 and Fig.3.1.3 show a comparison between the three models considered above, i.e. the non-degenerate ideal gas, the partially-degenerated Thomas-Fermi and the fully degenerate Fermi-Dirac.

**Fig.3.1.2.** Density profiles of the plasma produced by interaction of a 12 ps, 0.25 \( \mu \)m laser pulse with a solid Al target at an intensity of \( 5 \times 10^{16} \) W/cm\(^2\), 2.8 ps after the peak of the laser pulse. The dashed line was obtained modelling the electron gas with an equation of state based on the Fermi-Dirac distribution function while the solid and the dashed-dotted line were calculated assuming a perfect gas and a Thomas-Fermi model respectively.
The mass density of the plasma and the corresponding electron temperature are plotted as a function of the distance from the original target plane, at 2.8 ps after the peak of the 12 ps laser pulse. The laser intensity on target was \(5 \times 10^{16}\) W/cm\(^2\), i.e. close to the maximum laser intensity taken into account from an experimental viewpoint. One can see that the non-degenerate ideal gas EOS (dashed-dotted line) and the fully degenerate Fermi-Dirac EOS (dashed line) lead to identical results. As already discussed above, this is not surprising since degeneracy effects only occur in the cold supercritical plasma region, in the inner part of the target.

![Electron temperature profile](image)

**Fig.3.1.3**. Electron temperature profiles calculated using the same condition of Fig.3.1.2 at 2.8 ps after the peak of the laser pulse. A 12 ps, 0.25 µm laser pulse was set to interact with a solid Al target at an intensity of \(5 \times 10^{16}\) W/cm\(^2\).

In contrast, minor differences can be seen by comparing these curves (perfect gas and FD) with the analogous ones obtained assuming a TF model (solid line). In fact, according to Fig.3.1.1, in the density region relevant to these simulation results, the effects of electron-ion Coulomb interaction, taken into account by the TF model, are expected to be dominant over the quantum mechanical ones.

Nevertheless, these results indicate that, although non-ideal phenomena are clearly occurring in this regime, only minor effects can be seen on the hydrodynamic evolution of the plasma. It is however important to notice that the laser-matter interaction regime investigated in this section already accesses plasma conditions for which Coulomb interactions start playing a role in the thermodynamic properties of the plasma. The investigation of interaction regimes characterized by shorter laser pulses, higher intensities and shorter wavelength would necessarily have to fully include non-ideal plasma effects, as they are expected to play a crucial role.
3.2 - Coupling in Short Scalelength Plasmas

According to the theory summarized in Chpt.1 a plasma can support e.m.waves and, according to the dispersion relation given by Eq.1.2.10, propagation can take place as long as the angular frequency of the light is greater than the local plasma angular frequency, i.e. as long as the electron density is smaller that the critical density for the given laser frequency. Once laser light has entered the plasma, several mechanisms can lead to transfer of energy from the e.m.wave to the plasma. The two most important absorption mechanisms in the interaction regime typical of picosecond laser-produced plasmas are resonance absorption (RA) and collisional or inverse bremsstrahlung absorption (IB). The latter has already been introduced in Sect.1.2 and the absorption coefficient as a function of the plasma electron density and temperature is given by Eq.1.2.13. The effect of high laser intensity on this absorption coefficient will be discussed in Sect.3.3.

When the light propagates in an inhomogeneous plasma, whose scalelength is short compared to the typical collisional absorption length, resonance absorption can occur (Ginzburg, 1964) in the proximity of the critical density layer and up to 50% of the laser energy involved in the resonant process can be transferred to the plasma as electron plasma waves. In this section we will examine the role of collisional absorption and resonant absorption, during the propagation of laser light in short scalelength inhomogeneous plasmas, where both mechanisms can account for efficient laser-plasma coupling.

Inverse bremsstrahlung in inhomogeneous plasmas

As already pointed out, a necessary condition for resonant absorption to occur is that collisional absorption is not so efficient as to completely deplete the beam on its way to the critical density layer. It is therefore important to define the conditions for which RA can be expected to play an important role in the absorption processes in inhomogeneous plasmas. We will consider the general case of light incident on an inhomogeneous planar plasma, at an angle \( \theta \) with the plasma density gradient. The solution of the wave equation for obliquely incident light on an inhomogeneous plasma has been studied (Kruer, 1988), in the WKB approximation, for both linear
3.2 Coupling in Short Scalelength Plasmas

and exponential density profiles, taking into account the dependence of electron-ion collision frequency upon the electron density. According to the theory, the light propagates towards higher densities, up to the turning point, where it is reflected back. The density at which such reflection takes place is given by

\[ n_e = n_c \cos^2 \theta, \]  

(3.2.1)

where \( n_c \) is the critical density given by Eq.1.2.5. The fraction of the incident light absorbed by the plasma, due to collisional absorption, after reflection at the turning point is

\[ A = 1 - \exp \left( - \frac{b \nu_e^* L \cos^m \theta}{c} \right) \]  

(3.2.2)

where \( \nu_e^* \) is the electron-ion collision frequency given by Eq.1.2.12 and evaluated at the critical density, and \( L \) is the electron density scalelength. The coefficients \( b \) and \( m \) are respectively 32/15 and 5 in the case of a linear density profile \( n_e = n_c (1 - z/L) \), and 8/3 and 3 for an exponential density profile \( n_e = n_c \exp(-z/L) \). The exponential term of Eq.3.2.2, that is, the fraction of energy scattered by the plasma after reflection, has been plotted in Fig.3.2.1 as a function of the angle of incidence, for electron temperatures from 1 keV to 5 keV, for an exponential density profile of 4 µm scalelength and for 0.25 µm light. As will be shown in Chapt.5, these parameters are of particular relevance in the experimental study carried out on picosecond laser-plasma interaction.

**Fig.3.2.1.** Fractional scattering of laser light as given by Eq.3.2.1 for 0.25 µm laser light incident on a plasma with an exponential density profile with a 4 µm scalelength.
According to the plots of Fig.3.2.1, the fractional scattering increases (absorption decreases) with the increasing angle of incidence. This is mainly due to the fact that the distance of the turning point from the critical layer, where more efficient collisional absorption can take place, increases. In addition, the increase of the electron temperature makes the plasma less collisional, thus reducing the absorption. However, for small angles of incidence and for electron temperatures up to 2 keV, inverse bremsstrahlung accounts for absorption of most of the laser energy at 0.25 µm. These conclusions are based on the assumption that the absorption coefficient does not depend upon the laser intensity, i.e. linear IB absorption can be assumed. A discussion of non-linear absorption mechanisms occurring at high laser intensity will be given in Sect.3.3.

**Resonance absorption of P-polarized light**

This process consists of a linear conversion of the energy of the e.m. wave into longitudinal electron plasma waves. If a component of the laser electric field parallel to the density gradient exists when the light reaches the turning point, as in the case of P-polarized light, the electric field vector, oscillating at the frequency \( \omega_L \), will resonantly couple with the longitudinal electron plasma waves at the local frequency as given by Eq.1.2.2. Electron waves with angular frequency \( \omega_e = \omega_L \) will be excited and up to 50% of the e.m. wave energy can be converted into energy of the electron plasma waves. These waves will propagate down the density gradient and, due to collisionless Landau damping (see Eq.1.2.8) or collisional damping (see Eq.1.2.11), will loose their energy which will be eventually converted into hot electron energy (see Sect.1.2) or into thermal motion. According to the results shown above, short scalelength plasmas can offer suitable conditions for resonance absorption to occur, provided that the electric field vector of the laser light has a component in the plane of incidence.

According to Eq.3.2.1, as the angle of incidence decreases, the density at which the light is reflected approaches the critical density and, therefore, better resonance conditions can be achieved. However, the energy transfer from the light wave to the electron plasma wave also depends upon the time spent by the light wave in the region where resonant coupling can occur, which increases with the angle of incidence. Therefore, maximum energy transfer occurs at a particular angle which depends upon the density scalelength and the wave-number of the incident light. According to detailed numerical calculation (Forslund et al., 1975), the angle \( \theta_{\text{max}} \) of
maximum absorption for propagation of P-polarized light of wave-number $k$ incident onto a linear density gradient with scalelength $L$ is given by

$$\sin \theta_{\text{max}} \equiv \left(\frac{1}{2} k L\right)^{1/3}. \quad (3.2.3)$$

The angle $\theta_{\text{max}}$ obtained from Eq.3.2.3 has been plotted in Fig.3.2.2 as a function of the density scalelength for the wavelength of 0.25 $\mu$m considered above and for other two wavelengths of interest, i.e. the fundamental and second harmonic of the Neodymium glass laser. According to Fig.3.2.2, sizeable angles of peak absorption, say $\theta_{\text{max}} \geq 5$ deg, can be obtained if the density scalelength is smaller than 10 $\mu$m. For example, the peak angle for 0.25 $\mu$m light for the 4 $\mu$m scalelength plasma already considered above in this section is approximately 10 degrees. This value has in fact been measured experimentally as described in detail in Sect.5.2.

![Fig.3.2.2](image-url)  
Fig.3.2.2. Angle of maximum resonance absorption as a function of the scalelength for P-polarized light incident on an inhomogeneous plasma with a linear density profile.

The angular dependence of the fractional absorption due to the resonance process has been obtained with a simple calculation, based on a so called capacitor model (Mulser, 1991). According to this result the absorption coefficient $A$, as a function of the angle of incidence $\theta$, is given by the following implicit expression

$$A = 0.53 \left(2 - \sqrt{A}\right)^2 (kL)^{2/3} \frac{\sin^2 \theta \exp\left(-\frac{4}{3}kL \sin^3 \theta\right)}{\sqrt{1 - \sin^2 \theta}}, \quad (3.2.4)$$

where $k$ is the wave-number of the incident e.m.wave. Fig.3.2.3 shows the angular dependence of the RA coefficient obtained according to Eq.3.2.4 for the same conditions taken into account so far, that is for 0.25 $\mu$m P-polarized light incident
onto an inhomogeneous plasma with a scalelength of 4 µm. A correction factor has been introduced in order to match the maximum absorption with the prediction of the detailed theory (Forslund et al., 1975). This simple model reproduces the basic physical features of the resonance absorption. The absorption predicted by Eq. 3.2.4 peaks at an angle very similar to that given by Eq. 3.2.3. Also, consistently with the detailed theory, the range of angles of incidence around the peak Δθ, for which sizeable absorption occurs, is of the order of the peak angle itself, that is Δθ ≈ θ\text{max}.

![Absorption Coefficient vs Incidence Angle](image)

**Fig. 3.2.3.** Angular dependence of the resonance absorption coefficient for 0.25 µm P-polarized light incident on an plasma with a 4 µm scalelength linear density profile.

In the analysis of RA considered so far, a cold plasma approximation was assumed. On the other hand it has been shown (Speziale et al., 1977) that electron temperature effects can be expected only in a relativistic regime, where the electron thermal velocity approaches the speed of light. In this regime, the absorption coefficient approaches unity and the angle of maximum absorption increases. However, for electron temperatures up to several tens of keV there is almost no variation in the absorption coefficient, and the cold plasma approximation leads to an accurate description of the resonance absorption process.

**Hot electron generation**

As already discussed previously, once longitudinal electron plasma waves have been excited, collisional as well as collisionless damping processes lead to plasma heating. In the case of hot inhomogeneous plasmas, however, collisionless Landau damping is typically more efficient than collisional damping. In Sect. 1.2 it was
pointed out that Landau damping leads to the production of hot electrons, that is, to electrons with a kinetic energy much larger than the thermal energy. Detailed 2D simulations of the interaction of intense laser light with inhomogeneous plasmas have shown (Forslund et al., 1977) that the energy distribution function of these electrons is quasi-Maxwellian, with a characteristic temperature which scales approximately as $\Theta_h \propto (I\lambda_o^2)^{1/3}$. In addition, the fraction of hot electrons at the critical density is expected to be

$$\frac{n_h}{n_c} \approx \frac{A}{2F_h} \frac{\eta_h^2}{v_h}, \quad (3.2.5)$$

where $A$ is the absorbed fraction of laser energy, $F_h$ is the flux limiter associated to the hot electrons (see Eq.2.1.4), and $\eta_h = v_q/v_h$ is the ratio between the quiver velocity $v_q$ introduced in Sect.1.2, and the velocity associated to the characteristic hot electron temperature, $v_h$. This result indicates that, when the quiver velocity is of the order of the thermal velocity, the hot electrons density can be as high as 20% of the critical density.
3.3 - Non-Linear Energy Transfer Mechanisms

Non-linear effects in the interaction of laser light with plasmas take place when the quiver velocity of the electrons, in the oscillating laser electric field, approaches the electron thermal velocity. A simple way to characterize the degree of non-linearity in the interaction regime under investigation consists in evaluating the electric field strength parameter \( \eta = \frac{v_q}{v_{th}} \), i.e. the ratio between the electron quiver velocity \( v_q \) and the electron thermal velocity \( v_{th} \). When \( \eta \approx 1 \), the physics of the interaction discussed so far must be modified to include the dependence upon laser intensity. This section will be devoted to study the effect of laser intensity on the collisional absorption process discussed in Sect.1.2 and Sect.3.2.

**Langdon effect**

When the laser electric field interacts with the electrons in the plasma, energy is transferred from the e.m.wave to the electrons. The power density transferred to the electrons is given by the electron-ion collision frequency \( \nu_{ei} \), given by Eq.1.2.12, times the energy density of electrons oscillating in the laser electric field \( n_e m_e v_q^2 / 2 \). Since \( \nu_{ei} \) is inversely proportional to the third power of the electron velocity, lower velocity electrons will be preferentially heated, and the low velocity part of the electron distribution function will be depleted. However, electron-electron collisions can compensate this depletion, by redistributing the absorbed energy among all electrons. The power density redistributed by electron-electron collisions can be written as the e-e collision frequency, \( \nu_{ee} = \nu_{ei} / Z \), times the thermal kinetic energy density of electrons, \( n_e m_e v_{th}^2 / 2 \), \( Z \) being the plasma charge state. This process ensures that the Maxwellian distribution of electron velocities is preserved, provided that the equilibration rate is much greater than the rate of energy input. Therefore, the condition for Maxwellian electron distribution to be preserved can be written (Langdon, 1980)

\[
\frac{Z v_q^2}{v_{th}^2} \geq 1. \tag{3.3.1}
\]

When the external electric field is too high for the condition given by Eq.3.3.1 to be satisfied, electron-electron collisions are unable to preserve the Maxwellian distribution. The exact shape of the electron distribution function becomes a
function of the amplitude of the external driving field, and the heating process becomes non-linear. Since lower velocity electrons are preferentially heated by the e.m.wave, and redistribution does not take place, the expected depletion of the low velocity part of electron distribution (Langdon effect) will give rise to a super-Gaussian distribution. It has been shown that this depletion of the distribution function near $v = 0$ can account for a reduction of the inverse bremsstrahlung absorption coefficient of up to a factor of two, compared with the case of a Maxwellian distribution. The following correction to the linear coefficient has been suggested (Langdon, 1980) in order to take into account the Langdon effect

$$\kappa_{ib}^{Lang} = \kappa_{ib} \left(1 - \frac{0.553}{1 + \left(0.27/(Z\eta^2)\right)^{3/4}}\right),$$

(3.3.2)

where $\kappa_{ib}$ is the linear absorption coefficient given by Eq.1.2.13. The Langdon effect has in fact been observed in Fokker-Planck simulations of short pulse laser-plasma interactions (Rickard et al., 1989) even though the incident laser intensity was only $6 \times 10^{15}$ W/cm$^2$. Depletion of low-velocity electrons, characteristic of the Langdon effect, was observed as result of laser heating. As a final remark on the Langdon effect, we observe that the condition expressed by Eq.3.3.1 can be rewritten as $\eta \geq Z^{-1/2}$. Consequently, the higher the plasma charge state, the lower the laser intensity required for the Langdon effect to occur.

### Standard non-linear absorption

If the external electric field is so high that the electron quiver velocity is greater than the thermal velocity, i.e. $\eta = v_q/v_{th} > 1$, the electron distribution function will be entirely perturbed, the quivering motion being dominant over the thermal one. Although this process occurs at higher laser intensities than those required for the Langdon process, there is no limit to its effect on the absorption coefficient.

A simple but effective way to account for this non-linear effect consists in evaluating the inverse bremsstrahlung absorption coefficient given by Eq.1.2.13, replacing the electron temperature with an effective temperature, according to the equation $K_B T_e^{eff} = K_B T_e + m_e v_q^2$. The first term on the RHS of this equation is the average thermal kinetic energy while the second term accounts for the average kinetic energy of the quivering motion of the electrons.
Factorizing $T_e$ in the last equation, and expressing the electron temperature as a function of the thermal velocity, we obtain the following expression for the effective temperature as a function of the strength parameter, $\eta$

$$T_{e}^{\text{eff}} = T_e (1 + \eta^2),$$  \hspace{1cm} (3.3.3)

Substituting this expression into Eq.1.2.13, which gives the linear inverse bremsstrahlung absorption coefficient as a function of the electron temperature, we find the following result for the non-linear inverse bremsstrahlung coefficient

$$\kappa_{ib}^{NL} = \kappa_{ib} (1 + \eta^2)^{-3/2}.$$  \hspace{1cm} (3.3.4)

The collisional absorption depletion factor given by Eq.3.3.4, i.e. the ratio $\kappa_{ib}^{NL} / \kappa_{ib}$, has been plotted in Fig.3.3.1 as a function of the strength parameter. In the same graph, the depletion factor resulting from the Langdon effect, for a fully ionized Al plasma, has also been plotted for comparison.

![Fig.3.3.1](image)

**Fig.3.3.1.** Dependence of the collisional absorption reduction factor as a function of the electric field strength parameter calculated (Langdon, 1980) for the Langdon effect (dotted curve) for a fully ionized Al plasma and for the standard non-linear effect (solid curve) as given by Eq.3.3.4.

The reduction factor due to the Langdon effect, although important even for values of $\eta$ definitely smaller than unity, saturates to $\eta \approx 0.4$. In contrast, the standard non-linear effect becomes dominant for $\eta > 0.8$ and brings the absorption coefficient down to 10% of the Maxwellian value already at $\eta = 2$. Experimental measurements on absorption mechanisms will be described in Sect.5.2 where, according to the results presented here, the standard non-linear effect will be dominant over the Langdon effect.
Non-linear absorption in steep density gradients

In view of the application of these results to the case of propagation of laser light in steep density gradients, one should also take into account that, as discussed in Sect. 1.3, the swelling of the laser electric field occurring in the density gradient will further contribute to establish the conditions for non-linear processes to occur. In this case, the laser light propagates towards higher densities, up to the turning point where it is reflected, and then propagates backwards.

According to Eq. 1.3.1, the laser electric field increases as the density increases, and so does the parameter $\eta$ considered above. On the other hand, the collisional absorption coefficient also increases strongly with the electron density. Consequently, absorption will predominantly occur in the density region where strong swelling of the electric field occurs. Collisional absorption of laser light propagating in a plasma, in the non-linear regime, is generally described by the following differential equation

$$\frac{d I_L(x)}{dx} = -\kappa^{NL}_{\text{ib}}(x, I_L(x)) I_L(x),$$  \hspace{1cm} (3.3.5)

where $I_L(x)$ is the intensity of the laser light which has propagated for a distance $x$ in the plasma and the non-linear absorption coefficient $\kappa^{NL}_{\text{ib}}(x, I_L(x))$ is given by Eq. 3.3.4 and Eq. 1.2.13. The absorption coefficient depends upon the position, $x$ as well as upon the local value of the laser intensity, and contains the information on the given density profile. This first order differential equation can be solved for any density profile using numerical methods. Solutions of this equation, for linear as well as exponential density profiles, obtained with the Runge-Kutta method (Conte et al., 1981), will be discussed in Chapt. 5.

A simple preliminary estimate of the importance of non-linear effects on collisional absorption can be obtained evaluating the absorption coefficient as a function of the electron density, for propagation in a linear density profile. The non-linear coefficient given by Eq. 3.3.4 was evaluated for interaction conditions relevant to the experimental investigation presented in Chapt. 4, that is, for 0.25 µm laser light incident, at an intensity of $5 \times 10^{16}$ W/cm$^2$, on a 4 µm scalelength plasma at an electron temperature of 2 keV. Swelling of the electric field was also accounted for in the calculation. The result is shown in Fig. 3.3.2 where the IB absorption coefficient is plotted as a function of the electron density for both the linear and non-linear case. According to this graph, the standard non-linear effect strongly reduces the effectiveness of collisional absorption in this interaction regime.
For example, the absorption length at $0.85 \, n_e$ increases from the 3.57 µm of the linear case to 6.67 µm in the non-linear case. In these circumstances and in the presence of the critical density layer, one can expect that the resonance process described in Sect.3.2 will eventually play a key role in the absorption process. The correct conditions are consequently established for an interaction regime to take place, in which non-linear processes strongly affect the interplay between collisional and resonance absorption. This aspect will be further clarified in the discussion of the experimental results presented in Chapt.5.
3.4 - Limits of Steady-State Atomic Physics

In the general situation where plasma hydrodynamic conditions change on a time-scale fast compared to the typical time-scale of atomic processes, the population of excited and ionized states are not in equilibrium, therefore the steady state approximation is no longer valid, and the time-dependent general rate equations will have to be solved. In this section we will derive the constraints on plasma conditions to be fulfilled for the steady-state modelling to be satisfactory.

Transition rates of the main processes

If ionization/recombination and excitation/decay processes are both fast compared to the time-scale of plasma hydrodynamics, then a steady-state description of the plasma atomic physics will be acceptable. In the general case all the processes described in Sect.2.2 must be considered. However, according to Table 2.2.1 and considering that, in laser produced plasmas, radiation is usually weakly coupled to ions and electrons, radiative ionization processes can be neglected. Collisional ionization and auto-ionization are the two most efficient ionization processes. Equilibrium is established when the ionization processes are balanced by the corresponding recombination processes, i.e. three-body and dielectronic recombination.

The time-constant of the transient phase, which characterizes the response of atomic processes to sudden changes in the hydrodynamic plasma state, is given by the inverse of the sum of the transition rates. A semi-analytic expression (Lotz, 1969) for the collisional ionization rate, from the ground state \( g \) of an ion with charge \( Z \), to an ion with charge \( Z+1 \), is given by

\[
S(\Theta_e, Z, g) = 2.97 \times 10^{-6} \frac{\xi}{E_{\infty}^Z \sqrt{\Theta_e}} E_I(U) \quad \text{(cm}^3 / \text{s}) , \quad (3.4.1)
\]

where \( \xi \) is the number of electrons in the outer shell of the ion being ionized, \( E_{\infty}^Z \) is the ionization energy of the ion in eV, \( \Theta_e \) is the electron temperature in eV. In this equation \( U = E_{\infty}^Z / \Theta_e \), and \( E_I(U) \) is the exponential integral given by

\[
E_I(U) = \int_0^U (\exp(-x) / x) \, dx .
\]
3.4 Limits of Steady-State Atomic Physics

The rate of the inverse process, namely the three-body recombination, can be related to the direct process by the principle of detailed balance. The three-body recombination rate from the ground state \( g \) of an ion with charge \( Z + 1 \), to an ion with charge \( Z \), is therefore

\[
\alpha_{3b} = 1.66 \times 10^{-22} n_e g_Z \frac{g_{Z+1}}{g_{Z+1}^2} \exp\left(\frac{U}{\Theta_e}\right) S(\Theta_e, Z, g) \quad \text{(cm}^3/\text{s}) \tag{3.4.2}
\]

where \( n_e \) is the electron density in \( \text{cm}^{-3} \), \( g_Z \) and \( g_{Z+1} \) are the statistical weights of the ground levels of the two ionization stages. In the assumption of optically thin plasma we can neglect radiative ionization, but we must consider radiative free-bound recombination, as it can strongly contribute in balancing ionization. The rate for radiative recombination of ions with charge \( Z + 1 \) and ionization energy of the recombined ion \( E_\infty^Z \) (eV) is

\[
\alpha_{fb} = 5.20 \times 10^{-14} (Z + 1) U^{3/2} \exp(U) E_I(U) \quad \text{(cm}^3/\text{s}) \tag{3.4.3}
\]

where \( E_I(U) \) is the exponential integral given above. Eqs.3.4.1-3 can be used to determine if the ionization processes are fast enough to ensure that the conditions for a steady-state modelling of the ionization equilibrium are fulfilled.

From the point of view of excitation and de-excitation mechanisms in an optically thin plasma, collisional excitation will be balanced by collisional de-excitation and radiative spontaneous decay. A general expression (McWhirter, 1965) for the collisional excitation rate for a transition from a lower level \( l \) and an upper level \( u \), separated by an energy \( \Delta E_{u,l} \) (eV) is

\[
X(T_e, u, l) = 1.6 \times 10^{-5} \frac{\langle G(u, l) \rangle f_{u,l} \exp\left(-\frac{\Delta E_{u,l}}{\Theta_e}\right)}{\Delta E_{u,l} \Theta_e^{3/2}} \quad \text{(cm}^3/\text{s}) \tag{3.4.4}
\]

where \( \langle G(u, l) \rangle \) is the Gaunt factor averaged over the Maxwellian distribution of velocities of the electrons with temperature \( \Theta_e \), and \( f_{u,l} \) is the oscillator strength. The collisional de-excitation rate \( Y(\Theta_e, u, l) \) can again be obtained by the excitation rate using the principle of detailed balance and is given by

\[
Y(T_e, u, l) = 1.6 \times 10^{-5} \frac{\langle G(u, l) \rangle f_{u,l} g_l}{\Delta E_{u,l} \Theta_e^{3/2} g_u} \quad \text{(cm}^3/\text{s}) \tag{3.4.5}
\]

where \( T_e \) and \( \Delta E_{u,l} \) are again expressed in eV and \( g_u \) and \( g_l \) are the statistical weights of the upper and lower level respectively.
Finally, the rate of radiative decay is given by the Einstein coefficient $A_{u,l}$ for spontaneous emission

$$A_{u,l} = 4.3 \times 10^7 \frac{B_l}{g_u} f_{u,l} (\Delta E_{u,l})^2 \text{ (s}^{-1}) \text{.} \quad (3.4.6)$$

However, we observe that, although both ionization and excitation processes, and their inverse processes, play a role in determining the response time of the atomic processes to changes in the plasma conditions, a quantitative analysis of the transition rates shows that excitation and de-excitation processes are in general faster than ionization and recombinations processes. Therefore, excited states of a given ion state can be considered, in most cases, in equilibrium with the corresponding ground state (Hutchinson, 1987).

An estimate of the relaxation time of the ionization processes can be obtained assuming a two level system in which transitions take place from one level to the other, with the transition rates given by Eqs.(3.4.1-3). The solution of the rate equation is characterized by an initial transient phase, during which the population of the two levels relax to the equilibrium condition, with a time constant given approximately by

$$\tau = \frac{1}{n_e(S + \alpha_3 b + \alpha_f b)} \text{.} \quad (3.4.7)$$

As the experimental investigation presented in Chapt.4 and Chapt.5 was carried out on plasmas produced by laser irradiation of Aluminium and plastic targets, with and without Oxygen, the relaxation time of Eq.3.4.7 has been evaluated for these elements, at various ionization stages from Be-like to H-like ions.

**Transient ionization in Al plasmas**

Fig.3.4.1 shows the relaxation time for Aluminium ions as a function of the electron temperature and for a density of $10^2 \text{ cm}^{-3}$, which is the critical density for 1µm laser light. The density of interest here is the critical density since most of the X-ray radiation emitted by the plasma and used for plasma diagnostic purposes, originates from a region close to the critical layer. The wavelength corresponding to the fundamental frequency of the Neodymium laser ($\lambda=1 \mu m$) and to its fourth harmonic ($\lambda=0.25 \mu m$) have been considered in Fig.3.4.1 due to their relevance in experimental studies. The plots of Fig.3.4.1 show a dramatic increase in the
relaxation time from He-like to H-like Al, compared to ionization from Be-like to Li-like and from Li-like to He-like. This is mainly due to the large increase of the ionization energy of He-like ions involving the tightly bound K-shell electrons.

![Figure 3.4.1](image1)

**Fig. 3.4.1.** Relaxation times of ionization for Aluminium ions from Be-like to H-like as a function of the electron temperature for a density of $1 \times 10^{21}$ cm$^{-3}$.

According to this plot, the time taken by a plasma with an electron density of $10^{21}$ cm$^{-3}$, to achieve equilibrium between Li-like ion and He-like ions is of the order of several hundred picoseconds. **Fig.3.4.2** shows a plot analogous to that of Fig.3.4.1, but obtained at a density of $1.5 \times 10^{22}$ cm$^{-3}$, which is the critical density at the fourth harmonic ($\lambda = 0.25 \mu$m) of the Neodymium laser.

![Figure 3.4.2](image2)

**Fig. 3.4.2.** Relaxation times of ionization for Aluminium ions from Be-like to H-like as a function of the electron temperature for a density of $1.5 \times 10^{22}$ cm$^{-3}$.
Relaxation to He-like Al ions is now reduced to a few tens of picoseconds or less, mainly as a consequence of the scaling of the time constant with the electron density given by Eq.3.4.7. A steady state modelling is therefore expected to provide an accurate description of X-ray radiation emitted during plasma formation by nanosecond laser pulses at a wavelength of 0.25 µm. In contrast, in the lower density case of Fig.3.4.1, that is, in the case of 1 µm laser light, more restrictive conditions are established. According to Fig.3.4.1, He-like and H-like X-ray emission during nanosecond plasma formation should be considered in a transient regime. On the other hand, it should be noted that this analysis includes collisional ionization, three-body and radiative recombination only. Although these are the most important processes, other processes including, for example, charge-exchange recombination, can contribute to a faster relaxation.

Transient ionization in low-Z plasmas

In the case of Oxygen and Carbon plasmas, equilibrium is established on a picosecond time-scale or, as in the higher density case, on a sub-picosecond time-scale. Fig.3.4.3 and Fig.3.4.4 show the relaxation time from He-like to H-like ions for these two atomic species, as calculated according to Eq.3.4.7. for the two values of the electron density considered above. The corresponding curve relative to the Al plasma is also shown for comparison.

![Graph showing relaxation times of ionization for Carbon, Oxygen and Aluminium ions from He-like to H-like as a function of the electron temperature for a density of 1.0×10^{21} cm^{-3}](image)

**Fig.3.4.3.** Relaxation times of ionization for Carbon, Oxygen and Aluminium ions from He-like to H-like as a function of the electron temperature for a density of 1.0×10^{21} cm^{-3}.
Although a detailed analysis is required, for the particular experimental conditions under investigation from case to case, a general conclusion on the validity of a steady-state model can already be formulated on the basis of the results obtained so far and for the particular plasma species considered. The use of medium $Z$ elements, like Aluminium, sets a lower limit on the electron density at which processes involving He-like and H-like ionization states can be regarded as stationary on a given time scale. In the conditions of laser-produced plasmas using nanosecond pulses, this relaxation time can be as high as several hundred picoseconds. In the case of low $Z$ elements, ($Z<10$), plasmas with an electron density above $10^{21}$ cm$^{-3}$ can be considered stationary on a picosecond timescale. These circumstances will play a crucial role in the analysis of the time-resolved X-ray spectra presented in the following chapters.
CHAPTER 4

LONG-SCALELENGTH PLASMA STUDIES

An experimental investigation of long-scalelength plasmas will be described in this Chapter. Various techniques for the production of laser plasmas, for interaction studies relevant to inertial confinement fusion, will be discussed. Long-scalelength plasmas were produced and fully characterized in terms of electron density and temperature, by using several diagnostic techniques, including interferometry, X-ray time-resolved spectroscopy and X-ray time-resolved imaging. A detailed analysis of these measurements will be presented, followed by a comparison of the experimental results with the predictions of the numerical simulations presented in the previous Chapters. Finally, the results of a detailed experimental study on filamentation instability will be presented and discussed.
4.1 - Nanosecond Laser-Plasma Studies

The success of the inertial confinement fusion (ICF) scheme (Brueckner & Jorna, 1974) relies on the possibility of achieving a high degree of spatial uniformity in the compression of spherical pellets containing the nuclear fuel. Many competitive techniques have been proposed whose effectiveness in improving the uniformity of irradiation are presently under investigation world-wide. The role played by irradiation non-uniformities in the laser-plasma interaction physics is however not satisfactorily understood. Nevertheless it is now clear that they can seed the conditions for the development of various laser-driven plasma instabilities detrimental to ICF including self focusing and the filamentation instability.

The enhancement of laser intensity non-uniformities consequent to the onset of the filamentation instability in the sub-critical region will necessarily affect the uniformity of the ablation surface, seeding favourable conditions for the development of other instabilities including the Rayleigh-Taylor instability and the thermal instability. On the other hand, as discussed in Sect.2.3, other instabilities including the stimulated Brillouin scattering the stimulated Raman scattering and the two plasmon decay can be efficiently activated by a local increase of irradiance subsequent to the onset of filamentation and/or self focusing.

It is therefore clear that the interaction of the laser light with the underdense plasma surrounding the ICF pellet, where FI and SF are likely to occur, plays a key role in the depletion of the compression uniformity. Consequently, the experimental investigation of the physics of the interaction of laser light with large underdense plasmas is of crucial importance for the achievement of a control and eventually a suppression of these instabilities.

In the typical experimental configuration a long-scalelength test plasma is preformed by exploding a thin foil target. A delayed interaction laser beam is then focused onto this plasma and the physics of the interaction is investigated by means of an as wide as possible range of diagnostics. In order for the outcome of these interaction experiments to be correctly interpreted a detailed knowledge of the physical conditions of the preformed plasma is of fundamental importance. The attention has therefore recently been focused on the development of techniques for a reliable production and accurate characterization of the preformed test plasmas (Seka et al., 1992).
Long scalelength plasmas have been produced and extensively characterized in order to perform interaction studies under controlled conditions. The experimental configurations used for the production of underdense plasmas will be described in this section. Sect.4.2 will be devoted to the description of the interferometric technique used for electron density measurements. The limits of this technique and the method used to extract the phase shift map from the fringe pattern of the interferograms will be discussed in detail. The experimental density profiles obtained from the interferometry will be presented and discussed in Sect.4.3. The sensitivity of the interferometer to small scale density perturbations will also be analyzed in the same section. A description of the techniques employed for time-resolved X–ray imaging will be given in Sect.4.4 while time-resolved X-ray spectroscopy used for measurements of electron temperature and its temporal evolution will be presented in Sect.4.5. The results of an experimental investigation of the filamentation instability is presented and discussed in Sect.4.6.

Experimental approach

Many different schemes for production of long scalelength underdense plasmas have been suggested and adopted so far and their suitability for laser-plasma interaction experiments has been discussed in detail (Gizzi et al., 1994). Presently two main configurations are used whose features are considered to be suitable for interaction experiments relevant to direct-drive ICF, namely the so called line focus geometry (LFG) and the cylindrical geometry (CG). Fig.4.1.1 shows schematically the main features of the targets used in these configurations.

**Fig.4.1.1** Target configuration and interaction beam arrangement for the line focus geometry (LFG) and cylindrical geometry (CG) used for studies of laser interaction with long scalelength plasma relevant to direct drive inertial confinement fusion studies.
4.1 Nanosecond Laser-Plasma Studies

In the LFG, the plasma is produced by irradiation of a mass-limited stripe target and the interaction beam is set to propagate along the longitudinal axis of the target. The longitudinal plasma density scalelength can be directly controlled by varying the length of the target while the transverse density scalelength is set by the hydrodynamics of the target explosion and is typically comparable to the interaction beam focal spot size. The interaction beam is therefore subject to refraction effects which make the propagation particularly sensitive to slight misalignment of the interaction beam with respect to the longitudinal axis of the plasma.

This problem is overcome in cylindrical geometry where a mass-limited dot target is used to generate the preformed plasma, with the interaction beam set to propagate perpendicularly to the target plane. In this case the transverse density scalelength is of the order of the target diameter, typically $\approx 1\text{mm}$, and therefore refraction effects on the interaction beam are much smaller than in the LFG. On the other hand, in this case interaction takes place in an expanding plasma characterised by high flow velocities in the direction of the interaction beam, which complicates the analysis of the interaction physics. Table II summarizes the experimental parameters relative to these two configurations used in our experimental investigation. Both configurations have been studied experimentally and the plasmas produced have been extensively characterized in terms of spatial and temporal features of the main parameters including electron temperature and electron density.

<table>
<thead>
<tr>
<th>Target shape</th>
<th><strong>LINE FOCUS</strong></th>
<th><strong>CYLINDRICAL</strong></th>
</tr>
</thead>
<tbody>
<tr>
<td>Target dimension</td>
<td>500µm $\times$ 300µm</td>
<td>400µm (diameter)</td>
</tr>
<tr>
<td>Target thickness, material</td>
<td>700nm, Aluminium</td>
<td>500nm, Aluminium</td>
</tr>
<tr>
<td>Heating beams</td>
<td>2+2 (opposite)</td>
<td>2+2 (opposite)</td>
</tr>
<tr>
<td>Heating wavelength</td>
<td>0.526µm (Green)</td>
<td>1.053µm (IR)</td>
</tr>
<tr>
<td>Heating Focal spot</td>
<td>800µm $\times$ 400µm</td>
<td>600µm (diameter)</td>
</tr>
<tr>
<td>Heating intensity</td>
<td>$\approx 10^{14}$ W/cm²</td>
<td>$\approx 10^{14}$ W/cm²</td>
</tr>
</tbody>
</table>

Table 4.1.1. Laser and target configuration parameters relative to the two schemes employed in the experimental study of laser interaction with long scalelength plasmas. Typical values are indicated for target thickness and heating intensity which were varied in order to change plasma conditions at the time of interaction.

Characterization studies as well as laser-plasma interaction studies have been carried out on such plasmas by means of a wide range of diagnostics including, time-resolved visible spectroscopy, time-resolved X-ray spectroscopy, time resolved
X-ray imaging and visible interferometry. Characterization studies focused on the problem of reproducibility and uniformity of the test plasma, whereas interaction experiments have mainly concentrated on stimulated Brillouin scattering and filamentation instability.

**Experimental set-up in the cylindrical geometry (CG)**

The experimental investigation described here was mainly performed at the SERC Central Laser Facility. A schematic diagram of the experimental set-up for the cylindrical geometry case is shown in the layout of Fig. 4.1.2. Four 600 ps, 1.053 µm beams of the Vulcan laser were used to preform the plasma.

These *heating* beams were focused with f/10 optics in two opposed pairs and superimposed on target in a 600 µm spot providing a total irradiance of up to 7.5×10^{13} W/cm² on each side of the target. The target, consisting of 0.5 µm thick, 400µm diameter Aluminium dots coated onto a 0.1 µm plastic stripe, was placed 6 mm before the focal plane of the four heating beams, in the centre of the 600 µm diameter spot. Since the focal depth of each beam was approximately 1 mm, the target was basically located in the near field of these beams. Typical near field
Nanosecond Laser-Plasma Studies

intensity non uniformities had a scalelength of a twentieth of the beam size and were limited to less than 10% of the average intensity. The interaction pulse used in the investigation of laser-driven plasma instabilities was delayed by typically 2.5 ns with respect to the heating pulses and focused, using an f/15 optics, onto the preformed plasma along the main target symmetry axis, at an irradiance ranging from \(10^{13}\) to \(5 \times 10^{14}\) W/cm\(^2\). The focal spot diameter was typically 100 µm, i.e. much smaller than the plasma transverse scalelength to avoid refraction effects at the plasma boundary.

**Experimental set-up in line focus geometry (LFG)**

In the LFG the target consisted of 700 µm long, 300 µm wide, 700 nm thick Aluminium layer coated onto a 0.1 µm plastic substrate. A schematic overview of the experimental set-up is shown in Fig.4.1.3. The four heating beams of the Vulcan laser were frequency doubled and superimposed on target, two on each side. The focusing lens of each heating beam was tilted in order to generate a focal spot elongated along the longitudinal target axis, thus minimising losses of laser energy due to mismatching between the target cross-section and the focal spot.

**Fig.4.1.3.** Experimental set-up for laser-plasma interaction studies in line-focus geometry. The typical target consisted of an Aluminium stripe, several hundred microns thick, coated onto a thin plastic substrate.
As in the case of the CG, plasma conditions at the time of interaction where partially controlled by delaying the interaction beam with respect to the heating beams and/or by varying the heating laser irradiance. A preliminary estimate of the delay and the intensity necessary to achieve the required plasma conditions was formulated on the basis of the predictions of the numerical hydro-code MEDUSA (Christiansen et al., 1974) as discussed in detail in Sect.2.1. A semi-analytical hydrodynamic model (London & Rosen, 1986) was also used as a baseline to obtain simple scaling laws for the main plasma parameters.
4.2 - Interferometric Study of Plasma Density

Interferometric techniques can be used in order to measure the electron density distribution of a plasma as well as its temporal evolution. A fringe pattern is generated by the interferometer which gives the refractive index of the plasma integrated along the line of sight. Under special experimental conditions and with appropriate assumptions on the symmetry of the plasma, Abel inversion techniques allow the density distribution to be obtained with a temporal resolution given by the length of the laser pulse used as a probe beam. A detailed interferometric study of the plasma produced in the cylindrical geometry was carried out and the results will be presented and discussed in this section along with a full description of the technique used to analyse the interferograms, based on the Fourier transform.

Interferometer set up for phase shift measurements

A 100 ps (FWHM), 1 µm beam of the Vulcan laser was frequency doubled, delayed, and used as a probe beam for interferometric measurements parallel to the target plane as indicated in Fig.4.1.2. A modified (Benattar et al., 1979) Nomarski interferometer (Nomarski, 1955) was employed in order to measure the plasma induced phase shift on the probe beam.

A schematic description of the interferometer set-up is shown in Fig.4.1.2. This device basically consists (Willi, 1988) of a polarised light interferometer which produces, by means of a Wollaston prism, two separate orthogonally polarised images of the plasma surrounded by an unperturbed background. Interference between each of the two images and the unperturbed background of the other images is achieved by placing a polarizer before the film plane, oriented at 45 degrees with respect to the two axes of polarization of both images produced by the prism. The fringe pattern was recorded onto a commercial B&W film.

It should be noted that since a 0.53 µm wavelength probe pulse was used, which has the same wavelength as the second harmonic of the heating and interaction pulses, sources of second harmonic emission generated during the event could also be imaged on the same film. Therefore using a narrow-band interference filter, typically of 50Å of bandwidth, placed before the film, it was possible to obtain time-
integrated second harmonic (SH) images of the plasma interaction region at 90 degrees with respect to the heating/interaction beam axis as shown in Fig.4.1.2. The spatial resolution of the imaging system was measured to be better than 10 µm in the object plane.

**Experimental results**

Interferograms of the plasma were taken at various delays relative to the peak of the heating pulses in order to monitor the temporal evolution of the electron density. It was found that, under analogous intensity conditions, interferograms were highly reproducible shot by shot. Two representative interferograms will be analysed here which were obtained at approximately the same heating irradiance and with the interaction beam turned off. Fig.4.2.1 shows an interferogram taken 4.3 ns after the peak of the heating pulses at a heating irradiance of $4.2 \times 10^{13}$ W/cm$^2$ on each side of the target. Only one of the two fringe systems produced by the interferometer is entirely shown in this figure. Fringes on the extreme right hand side of the image are perturbed due to the overlapping with the other fringe system. The original position of the target is marked by two arrows and the edge of the target holder is just outside the image.

![Interferogram of the preformed plasma taken 4.3 ns after the peak of the heating laser pulses. The heating irradiance on each side of the dot Al target was $4.2 \times 10^{13}$ W/cm$^2$.](image)

Moreover, a minor fraction of the energy in the tail of the focal spot of the heating beams hit the holder producing a tenuous plasma giving rise to a slight perturbation of the fringe pattern visible at the top and bottom edges of the image.
Such perturbations are however limited to marginal regions and are expected to give a slight contribution to the fringe pattern in the region close to the target.

The time, relative to the heating pulses, at which the interferogram of Fig. 4.2.1 was taken, was the earliest time at which fringes were visible over the whole plasma extent. As will be discussed in detail in Sect. 4.4, a depletion of the fringe visibility in the region around the peak density was observed at earlier times as shown in the next figure. Fig. 4.2.2 shows another interferogram of the preformed plasma taken 3.0 ns after the peak of the heating pulses, at a heating irradiance of $3.2 \times 10^{13}$ W/cm$^2$ on each side of the target.

Both interferograms show that, away from the target plane, the plasma extends up to distances from the longitudinal axis greater than the target radius. These circumstances suggest that two dimensional effects in the hydrodynamic motion of the plasma may play an important role. Moreover, the absence of strong perturbations of the fringe pattern on a small spatial scale indicates that, in the limit of sensitivity of our interferometer, the plasma can be considered free from strong density perturbations. A comparison between the experimental longitudinal density profiles and the ones calculated by the 1D Medusa hydro-code will be presented in Sect. 4.3 along with a discussion of the sensitivity of the interferometer to small scalelength density perturbations.
Plasma induced refraction effects

It should be noted that density measurements based on interferometric techniques rely on the assumption that the electron density is well below the critical value, and that transverse density gradients are small enough so that the probe beam can propagate through the whole plasma without undergoing severe deflection. In fact, strong refraction effects might destroy the coherence across the wave-front of the probe beam. A simple estimate of the validity of the last assumption, which is also the most restrictive one, has been obtained in the approximation of a cylindrical plasma with a parabolic density profile.

In this case the maximum angular deflection is simply given by the peak density divided by the critical density. The phase disruption induced by bending effects is negligible provided that (Hutchinson, 1987) the maximum deflection angle is much smaller than the acceptance angle of the interferometer \( \alpha \), that is, \( n_e/n_c @ 0.5 \mu m << \alpha \), where \( n_e \) is the peak electron density and \( n_c @ 0.5 \mu m \) is the critical density at the probe wavelength. According to the simulation, 2.0 ns after the peak of the heating pulses, the peak density on the target plane is approximately \( 1.7 \times 10^{20} \text{cm}^{-3} \), that is, \( n_e/n_c @ 0.5 \mu m \equiv 0.05 \). Since the interferometer had an \( f/4 \) optics which corresponds to an acceptance angle of 0.25 rad, we can conclude that in our experimental conditions refraction effects are not expected to significantly perturb the transverse coherence of the probe beam.

Basic principles of interferometric analysis

The fringe pattern produced by the interferometer is the result of interference between a beam which has propagated through the plasma and an unperturbed reference beam, both beams originating from the same laser source. If the condition stated above is fulfilled, the probe beam propagates through the plasma in a straight line. According to the reference frame of the experimental set up of Fig.4.1.2, the phase difference between these two beams at a given position \((x, z)\) in an output plane perpendicular to the probe beam is

\[
\Delta \varphi (x, z) = \frac{2\pi}{\lambda_p} \left\{ \int_{-L/2}^{L/2} (\varepsilon(x, y, z) - 1) \, dy \right\}, \quad (4.2.1)
\]

where \( \lambda_p \) is the probe beam wavelength, \( \varepsilon \) is the plasma refractive index and \( L \) is the total path-length greater than the plasma extent along \( y \). According to Eq.1.2.6 and considering that \( n_e \ll n_c \), the plasma refractive index is given by
4.2 Interferometric Study of Plasma Density

\[ \epsilon(x, y, z) \equiv 1 - \frac{n_e(x, y, z)}{2n_{c@0.5\mu m}}, \quad (4.2.2) \]

where, according to Eq.1.2.2, the plasma frequency has been expressed in terms of
the electron density, \( n_e(x, y, z) \) and, according to Eq.1.2.5, the probe wavelength has
been expressed in terms of the corresponding critical density, \( n_{c@0.5\mu m} \). Eq.4.2.1
therefore defines a relation between the density distribution and the phase shift
distribution on the output plane of the interferometer.

The determination of \( n_e(x, y, z) \) from the phase shift distribution or, in other
words, the inversion of Eq.4.2.1 is, in general, a very difficult task, unless the plasma
has a cylindrical or spherical symmetry. Since in our case both the target and the
focal spots of each heating beam satisfy cylindrical symmetry around the \( x \)-axis,
one can expect that the plasma itself will show the same symmetry.

![Fig.4.2.3](image)

**Fig.4.2.3** A plasma with a cylindrical symmetry can be studied as a 2D problem by
considering a plane perpendicular to the symmetry axis. The reference frame used here is consistent
with the one shown in the experimental set up. (see Fig.4.1.2)

If the plasma has a cylindrical symmetry around the \( x \)-axis, with the probe beam
propagating along the \( y \)-axis, then the problem can be treated as two dimensional by
considering a plane perpendicular to the \( x \)-axis, say \( x = x_o \). According to the
definitions of Fig.4.2.3, Eq.4.2.1 at \( x = x_o \) becomes

\[ \Delta \varphi_{x_o}(z) = \frac{2\pi}{\lambda_o} \left( \int_L \left( \epsilon(r) - 1 \right) \frac{rdr}{\sqrt{r^2 + z^2}} \right), \quad (4.2.3) \]

where the refractive index of the plasma now depends upon \( r \), the distance from the
cylindrical symmetry axis. Substituting Eq.4.2.2 for the plasma refractive index in
Eq.4.2.3 and expressing, according to Eq.1.2.5, the critical density in terms of the
wavelength of the probe beam, one obtains the phase difference as a function of \( z \),
for a given plane \( x = x_o \), in the form of the Abel integral equation
\[ \Delta \varphi_{x_0}(z) = -\frac{e^2 \lambda_p}{m_e c^2} \int z \frac{n_e(r)}{\sqrt{r^2 - z^2}} dr, \]  

(4.2.4)

where \( e \) and \( m_e \) are the electron charge and mass respectively and \( c \) is the speed of light and \( r_o \) is the radial plasma size. Inversion of this equation leads to the solution for \( n_e(r) \)

\[ n_e(r, x = x_0) = -\frac{m_e c^2}{\pi e^2 \lambda_p} \int r \left( \frac{d}{dz} \Delta \varphi_{x_0}(z) \right) \frac{dz}{\sqrt{z^2 - r^2}}, \]  

(4.2.5)

in terms of the first derivative of the phase difference along \( z \) at \( x = x_0 \). According to this equation one can obtain the electron density once the phase shift induced by the plasma has been measured.

In order to detect this phase shift a small angle is introduced by the interferometer between the probe beam and the reference beam so that, in the absence of plasma, a pattern of parallel fringes is generated, as shown on the extreme left of Figs.4.2.1-2. The phase shift introduced by the plasma is then evaluated by measuring the displacement of the fringes from their unperturbed position. By counting the number of fringes crossed moving along \( z \) for a given \( x = x_0 \) one can obtain a sampling of \( \Delta \varphi_{x_0} \). This set of data can be fitted and introduced in Eq.4.2.5 in order to calculate the electron density.

However this procedure, often used in the past, is subject to a number of uncertainties introduced by the small amount of data available to build \( \Delta \varphi_{x_0} \) and by the particular choice of the fitting function. Moreover, the sensitivity of this technique is basically limited to one fringe shift which, as will be clear in the next section, would strongly limit the effectiveness of the technique itself, capable of a much better resolution. The use of a Fourier transform based method in the analysis of the interferograms allows a much more direct approach to the problem, free from the uncertainties described above, and easy to implement with simple numerical techniques as discussed below.

**Fourier analysis of the interferograms.**

The intensity of the fringe pattern produced by the interferometer on the film in the presence of plasma can be written as

\[ I(x, z) = a(x, z) + b(x, z) \cos[2\pi f_o x + \Delta \varphi(x, z)] \]  

(4.2.6)

where a 1:1 ratio of plasma to image magnification has been assumed. In this equation \( a(x, z) \) and \( b(x, z) \) account for non-uniformities in the background.
4.2 Interferometric Study of Plasma Density

Intensity and fringe visibility, $f_u$ is the spatial frequency of the unperturbed fringe pattern, i.e. the number of fringes per unit length on the film, $\Delta \varphi(x, z)$ is the phase shift given by Eq. 4.2.1. It has been shown (Takeda et al., 1982) that, by means of Fourier analysis, the intensity $I(x, z)$ can be processed in order to directly obtain the phase shift $\Delta \varphi(x, z)$. By expressing the cosine function of Eq. 4.2.6 in terms of the exponential function, Eq. 4.2.6 becomes

$$I(x, z) = a(x, z) + [c(x, z) \exp(2\pi i f_u x) + c. c.], \quad (4.2.7)$$

where $c(x, z) = (1/2)b(x, z)\exp[i \Delta \varphi(x, z)]$ and its complex conjugate $c^*(x, z)$ carry all the information concerning $\Delta \varphi(x, z)$. According to the definition of the logarithm of a complex number one can verify that

$$\log[c(x, z)] = \log[(1/2)b(x, z)] + i \Delta \varphi(x, z). \quad (4.2.8)$$

The phase shift can therefore be obtained by taking the imaginary part of the complex logarithm of $c(x, z)$. On the other hand, if we take the Fourier transform of Eq. 4.2.7 with respect to $x$ we obtain

$$F_i(f, z) = F_a(f, z) + F_c(f - f_u, z) + F_c^*(f + f_u, z), \quad (4.2.9)$$

where $F_a(f, z)$ is the Fourier transform of the background intensity along $x$ and $F_c(f - f_u, z)$ and $F_c^*(f + f_u, z)$ are the Fourier transforms of the two terms containing $\Delta \varphi(x, z)$. If the scalelength of typical non-uniformities of the background intensity along $x$ is large compared to the fringe frequency, then the contribution of $F_c(f, z)$, to the total Fourier spectrum of Eq. 4.2.9, will be well separated from the contribution due to the background intensity non-uniformities.

In this case, we can extract $F_c(f - f_u, z)$ (or $F_c^*(f + f_u, z)$) from the spectrum, shift it by $f_u$ along the frequency axis toward the origin in order to obtain $F_c(f, z)$, and perform the inverse Fourier transform to obtain $c(x, z)$. According to Eq. 4.2.8, the imaginary part of the complex logarithm of $c(x, z)$ will give the phase distribution $\Delta \varphi(x, z)$, that is still locally indeterminate by a factor of $2\pi$ resulting from the use of an inverse trigonometric function to obtain the argument of $c(x, z)$. However, this indetermination can be solved by setting an appropriate algorithm able to detect and compensate for jumps in the phase shift.

The interferograms were digitised with the two scanning directions set along $x$ and $z$ respectively (see Figs. 4.2.1-2). The optical density of the film was converted into intensity and stored in a two-dimensional array. A fast Fourier transform (FFT) of the intensity distribution along the direction perpendicular to the fringes, i.e.
along $x$, was performed for each position along $z$. **Fig.4.2.4** shows the modulus of the Fourier transform of the interferogram of Fig.4.2.1, with the origin of the spatial frequency axis set to be in the centre of the graph.

The two side components symmetric to the zero frequency correspond to the fringe pattern while the central, strong component accounts for the low spatial frequency variations of the background intensity. A portion of the interferogram larger than that shown in Fig.4.2.1, and containing more unperturbed fringes, was used to generate the pattern shown in Fig.4.2.4 in order to obtain a greater accuracy in the determination of the frequency of the unperturbed fringes. In fact, according to the shift theorem of Fourier theory, an error in the determination of the frequency shift, if applied to one of the side components as discussed above, would result (Nugent, 1985), after Fourier inversion, in a modulation of the final phase distribution along $x$ with a spatial frequency equal to the error itself. On the other hand, the finite sampling interval of the frequency space resulting from the finite sampling interval of the interferogram itself sets a lower limit to the uncertainty in the determination of the required frequency shift.

**Fig.4.2.5** shows a 3D shaded surface of the phase shift distribution generated by the preformed plasma 4.3 ns after the peak of the heating pulses, as obtained from the interferogram of Fig.4.2.1, while **Fig.4.2.6** shows the corresponding contour plot with a contour interval of $\pi$. 

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**Fig.4.2.4.** Fast Fourier transform of the intensity profiles of the interferogram of Fig.4.2.1. The FFT was performed along the direction perpendicular to the unperturbed fringes. The natural logarithm of the modulus of the Fourier transform is shown here as a grey-scale density plot.
Fig. 4.2.5. Three dimensional shaded surface of the phase shift distribution induced by the preformed plasma, 4.3 ns after the peak of the heating laser pulses. The phase shift was obtained from the interferogram of Fig. 4.2.1 using a fast Fourier transform based analysis technique.

Fig. 4.2.6. Contour plot of the phase shift distribution of Fig. 4.2.5 with a contour interval of \( \pi \), the uppermost curve corresponding to \( \varphi = \pi \).

Fig. 4.2.6 shows that the contour lines are highly symmetric with respect to the \( x \)-axis in the \( x-z \) plane. This is consistent with the assumption of cylindrical
symmetry of the plasma around the $x$-axis. An analogous technique has been employed in order to obtain the phase shift distribution from the interferogram of Fig.4.2.2, relative to the preformed plasma, 3.0 ns after the peak of the heating pulses. A contour plot, again with a contour interval of $\pi$, of the side with $z > 0$ is shown in Fig.4.2.7, the plot being symmetric with respect to the $x$-axis.

![Contour plot](image)

**Fig.4.2.7.** Contour plot of the phase shift distribution produced by the preformed plasma, 3.0 ns after the peak of the heating pulses. The phase shift distribution was obtained from the interferogram of Fig.4.2.2 using a technique based on the fast Fourier transform (see text).

The shaded area denotes the region where the phase shift could not be determined, the fringe visibility being below the detection level (see Sect.4.3).

The phase shift distribution shown in Figs.4.2.6&7 contains all the information necessary, once appropriate assumptions on plasma symmetry have been adopted, to determine the electron density distribution. These aspects will be discussed in the next section where density profiles obtained by Abel inversion are presented and compared with 1D numerical simulations.
4.3 - Electron Density Profiles

The phase shift distribution shown in Sect.4.2 can be treated mathematically in order to obtain the electron density distribution of the plasma that produced it. This procedure will be described in this section and the limits of the technique will be discussed, with particular attention to the sensitivity of the diagnostic, as a whole, to small scalelength density perturbations. A detailed analysis of the results and a comparison with 1D hydrodynamic simulations will also be presented.

From phase shift to density profiles.

According to Eq.4.2.5, the phase distributions of Fig.4.2.6 and Fig.4.2.7 can be inverted to determine the electron density distribution (Gizzi et al., 1994). The integral in the equation was solved numerically using a corrected composite trapezoid rule (Conte et al., 1980). Due to the large amount of data available, typically $512 \times 512$, the integral could be performed directly on the data itself, without polynomial fitting. Fig.4.3.1 shows a contour plot of the density profile obtained from Abel inversion of the phase shift of Fig.4.2.6.

![Contour plot of the electron density profile of the plasma 4.3 ns after the peak of the heating pulses as obtained from Abel inversion of the phase shift distribution of Fig.4.2.6. Contour levels are labelled in terms of the critical density at 1.053 µm.](image)

Fig.4.3.1
The contour levels in Fig.4.3.1 are labelled in terms of the critical density at 1 µm. Moreover, consistent with the assumption of cylindrical symmetry, the \( z \)-co-ordinate of Fig.4.2.5 has been replaced by the radial co-ordinate.

As already discussed in the description of the interferograms of Figs.4.2.1-2, some contribution to the phase shift comes from the plasma blow-off from the target holder. Since this plasma is basically confined into marginal regions, away from the central bulk of the plasma, its contribution to the phase shift was removed by multiplying the phase shift by an appropriate hanning function

\[
\Delta \varphi(x, z) \Rightarrow \Delta \varphi(x, z) \exp\left[-(z / w_z)^{2n}\right]
\]

where \( w_z \) and \( n \) were chosen in order to obtain a smooth cut-off of the unwanted contribution without modifying the phase shift due to the main plasma. Although this procedure introduces arbitrariness in the phase distribution of these marginal regions, the final density profiles obtained from Abel inversion were found to be rather insensitive to the particular choice of the values given to \( w_z \) and \( n \).

A similar procedure was adopted for the phase distribution of Fig.4.2.7 generated from the interferogram of Fig.4.2.2 which was taken 3.0 ns after the peak of the heating pulses. The lack of knowledge in the shaded region of the phase shift distribution of Fig.4.2.7 resulted in a corresponding indeterminacy in the electron density. Fig.4.3.2 shows a contour plot of the density profile obtained from Abel inversion of the phase shift distribution of Fig.4.2.7.

Fig.4.3.2 Contour plot of the electron density profiles of the plasma 3.0 ns after the peak of the heating pulses as obtained from Abel inversion of the phase shift distribution of Fig.4.2.7. Contour levels are labelled in terms of the critical density at 1.053 µm.
The contour lines for radial distances less than 100 µm indicate the presence of a minimum of the electron density on the target axis. This is likely to be an effect of the slightly higher laser intensity in the centre of the spot.

As already pointed out in the analysis of the contour plot of Fig.4.2.7, the shaded portion of the graph indicates the region where Abel inversion was not possible, as no interference fringes are visible in this region. Several processes have been taken into account in order to explain the observed effect and a detailed discussion is given below. The probe beam propagates through the plasma and therefore is subject to both absorption and refraction effects which can reduce its intensity and affect its coherence. Collisional absorption has been evaluated assuming propagation of the probe beam in a density profile similar to that of Fig.4.3.2. The peak electron density value was chosen according to the simulation, as given by Fig.2.1.3, and the electron temperature was set according to the results obtained from experimental measurements (see Sect.4.4). It was found that less than 10% of the probe energy would be absorbed by the plasma and consequently absorption is expected to have a marginal effect of the probe beam.

**Fringe visibility and density evolution.**

On the other hand, as the plasma expands, the density in the centre of the plasma decreases rapidly resulting in motion of the fringes on the image plane of the interferometer. Since the interferogram is recorded integrating over the whole probe pulse duration, if the rate of density change is such that fringes move by half of the fringe separation or more during this time, the fringe visibility will consequently vanish. For a simple estimate of this effect we take into account the result of the numerical hydrodynamic simulation that, in the proximity of the target plane, it is expected to be acceptable, as shown by Fig.4.3.4.

According to the simulation, 3 ns after the peak of the heating pulses, the electron density on the target plane decreases with a rate of $3 \times 10^{19} \text{ cm}^{-3}\text{ns}^{-1}$. Consequently, during the 100 ps probe pulse, the electron density varies by typically $3 \times 10^{18} \text{ cm}^{-3}$. This variation would lead to a phase shift of approximately $\Delta \varphi = 2.5 \text{ rad}$ which corresponds to almost half a fringe separation. Consequently the rate of density variation in the centre of the plasma could account for most of the loss of fringe visibility observed in Fig.4.2.2. In addition, we observe that, in agreement with experimental results, such an effect would give its maximum contribution to the depletion of fringe visibility in the centre of the plasma, where the rate at which the electron density decreases is maximum.
On the other hand, two dimensional effects in plasma expansion, if present, will give rise to a more rapid decrease of the electron density as the plasma will expand radially as well as longitudinally. Therefore the phase shift obtained above should be considered as a lower limit. All these effects tend to vanish at later times when the rate decreases. For example, 4.4 ns after the peak of heating pulses, the simulation predicts a rate of density change of $10^{19}$ cm$^{-3}$ ns$^{-1}$, resulting in a fringe shift during the probe pulse of $\Delta \varphi = 0.8$ rad, which corresponds to one eighth of the fringe separation. This is consistent with the interferogram of Fig.4.2.1, that shows fringes across the whole plasma extent.

**Small scale density inhomogeneities**

Since the fringe shift recorded on film is the result of spatial integration along the line of sight of the interferometer, the sensitivity of our measurements to local density inhomogeneities depends upon their scalelength as well as their amplitude. According to Eq.4.2.1 and Eq.4.2.2 the phase shift induced by the plasma depends linearly on the density itself. The contribution of a local density perturbation to the total phase shift can therefore be written as

$$
\delta \varphi(x, z) = - \frac{\pi}{\lambda \rho n_c @ 0.5 \mu m} \int_{-L/2}^{L/2} \delta n(x, y, z) \, dy,
$$

(4.3.2)

where $\delta n(x, y, z)$ describes the density inhomogeneity. In other words, in the limit of validity of the approximation of Eq.4.2.2, the phase shift induced by the density fluctuation is independent from the background density. We evaluated $\delta \varphi$ assuming a density inhomogeneity along the line of sight at a given position $(x_o, z_o)$ on the output plane of the interferometer given by

$$
\delta n_c(x, y, z) = a n_c @ 1 \mu m \exp\left(-y^2/w^2\right),
$$

(4.3.3)

where $a$ is the amplitude of the density perturbation in units of the critical density at 1 \(\mu\)m, $n_c @ 1 \mu m$ and $w$ is the scalelength of the perturbation. Fig.4.3.3 shows a contour plot of $N = \delta \varphi(0, 0)/2\pi$, that is, the number of fringe shifts corresponding to the phase shift, as a function of the two parameters $a$ and $w$, with $a$ ranging from zero to 0.05 and $w$ ranging from zero to 50 \(\mu\)m. The range of perturbation scalelength considered in Fig.4.3.3 is of particular relevance to laser driven instabilities in coronal plasmas as discussed in Sect.4.6. In this region of
parameters, the perturbation to the phase shift is expected to be less than $2\pi$ and therefore less than a fringe shift on the interferogram.

On the other hand, the use of the Fourier technique for the analysis of interferograms allows a very accurate determination of the phase shift distribution. The limit of this accuracy can be estimated from the contour plot of Fig.4.2.6. and is basically given by the small scalelength noise which affects the position of the contour curves. This effect leads to an uncertainty on the phase shift of the order of 0.5 rad, that is, less than one tenth of a fringe separation.

![Fig.4.3.3 Contour plot of the phase shift (in number of fringe shifts) induced by a Gaussian perturbation in the electron density along the line of sight of the interferometer as a function of the transverse perturbation scalelength, $w$, and the perturbation amplitude, $a$, in units of the critical density at a wavelength of 1.053 µm.](image)

Consequently, a single density perturbation along the line of sight with a scalelength of 20 µm would be detected as long as the corresponding amplitude of density fluctuation is greater than $0.01 n_{e@\lambda_{1\mu m}}$. This limit becomes $0.025 n_{e@\lambda_{1\mu m}}$ for a 10 µm scalelength perturbation. These conclusions make the interferometric measurements used in this work able to detect plasma inhomogeneities distinctive of most of the instabilities relevant to ICF presently under investigation.

The interferograms of the plasmas obtained in this experimental investigation generally showed a good small scale homogeneity, as shown by the two examples in Figs.4.3.1-2. Notice that, in Fig.4.3.1, the modulations visible on the contour line at the electron densities of $0.003 n_{e@\lambda_{1\mu m}}$ are beyond the sensitivity of the interferometer and consequently have no clear physical meaning. They are probably
related to the residual modulation of the phase distribution arising from the numerical analysis of the interferogram discussed in Sect.4.2.

On the other hand, local shift of the fringes corresponding to small scalelength density perturbations, induced by the delayed interaction pulse focused on the plasma, were indeed detected and will be discussed in Sect.4.6. Incidentally we observe that Eq.4.3.2 can also be used to determine the lowest electron density which can be detected for a given plasma extent along the line of sight of the interferometer. Assuming a plasma extent of the order of the target diameter, that is, 400 µm, the uncertainty of 0.5 rad on the phase distribution mentioned above gives a lower limit to the detectable density of \( \approx 10^{-3} n_e \, \mu \text{m} \) which is consistent with the measured limit given by Fig.4.3.1.

The evaluation of the sensitivity of the interferometer to small scale density non-uniformities presented above is a fundamental step for a correct interpretation of interaction experiments. From the point of view of the filamentation instability, for example, the conclusions drawn above indicate that the bulk of the plasma produced with our method is essentially free from those density perturbations able to efficiently initiate the instability. In fact, as discussed in Sect.4.6, a few nanoseconds after the plasma formation, when plasma conditions are suitable for ICF related interaction experiments, the optimum size for the growth of the filamentation instability is expected to be approximately 10-20 µm. On the other hand the interferograms indicate that the density perturbations in the bulk of the plasma in this range of scalelengths are \( \delta n_e / n_e < 0.2 \).

These conclusions cannot be extended to the low density plasma blow-off, the amplitude of density inhomogeneities in this region being below the detection level of the interferometer. Nevertheless, the plasma in this region will benefit, in terms of homogeneity, from hydrodynamic expansion processes which, at sufficiently large distances from the target plane, should contribute to smoothing out residual non-uniformities. However, from the point of view of laser-driven instabilities, the large density scalelength in the direction of propagation of the interaction beam establishes, by itself, favourable conditions for instabilities to grow.

**Effect of thermal smoothing on density inhomogeneities**

Electron density inhomogeneities can arise (Campbell, 1992) from intensity non-uniformities in the focal spot of the laser beam heating the target. The use of the fundamental harmonic of the Nd laser to preform the plasma, in place of its most commonly used second and third harmonic, has two important advantages. Firstly
laser intensity non-uniformities are those intrinsic to the laser system and, in the experiments presented here, were typically limited to 10% of the average intensity, with a scalelength of a twentieth of the beam size, i.e. approximately 30 µm on the target plane. Secondly, the use of longer wavelength light makes thermal smoothing more effective.

In fact the separation between the critical density layer and the ablation layer, during the initial explosion of the target, increases with the wavelength of the incident laser light. It has been shown (Key, 1991) that laser intensity fluctuations $\Delta I_L$, at the critical density layer, give rise to fluctuations in the thermal electron flux at the ablation layer $\Delta Q$, which depend upon the perturbation scalelength $L$ (cm), the average absorbed intensity $I_L (10^{13} \text{W/cm}^2)$ and the wavelength $\lambda_L (\mu\text{m})$ according to

$$\frac{\Delta Q}{Q} = \exp(-4.5 \times 10^7 I_L^{1/3} \lambda_L^{2/3} t L^{-1}) \frac{\Delta I_L}{I_L},$$  \hspace{1cm} (4.3.4)$$

where $t$ is the time from the beginning of irradiation in seconds. According to this relation, 100 ps after the beginning of laser irradiation at an average absorbed intensity of $10^{13} \text{W/cm}^2$, an intensity perturbation with a scalelength of 10 µm will give rise to a perturbation in the electron flux with the same scalelength at the ablation layer almost 100 times less intense. In the case of irradiation at the fourth harmonic of Nd laser ($\lambda_L = 0.25 \mu\text{m}$) the amplitude of the perturbation would be reduced only by a factor of 6. In other words, the laser wavelength sets an upper limit to the perturbation scalelength that will be efficiently smoothed out.

Thermal smoothing will also be effective in reducing residual non-uniformities after target explosion. In the plasma conditions of our experiment, at the end of the heating pulses, i.e. at $n_e = 10^{20} \text{cm}^{-3}$, $\Theta_e = 500\text{eV}$, the electron mean free path is a few microns. Therefore, plasma non-uniformities on a scale of less than $\approx 10 \mu\text{m}$ will also be smoothed out efficiently by thermal conduction after a few nanoseconds. There is recent evidence however, that in particular experimental conditions, thermal smoothing is less efficient than expected according to the current models. In fact it was found (Desselberger et al., 1992) that when plasmas are generated by irradiating solid targets, plasma density inhomogeneities, initiated early in the interaction by laser intensity non-uniformities, can persist during and after the laser pulse even in the presence of laser beam smoothing. Within the limits of sensitivity of the diagnostics employed in our experiment there is no evidence of similar processes occurring in our experimental conditions. In addition, the laser irradiance in our case has been limited to levels lower than those used for visible or
UV light, in order to avoid undesired non-linear effects which might destroy the plasma homogeneity. In fact these effects would be more important with 1 µm laser light as they scale roughly as $I_2 \lambda_\text{L}^2$. In the case of 1 µm wavelength and according to the intensity thresholds of the main non-linear processes active in laser-plasma interaction summarized in Sect.2.3, an intensity of $10^{14}$ W/cm$^2$ can be regarded as a cautious limit below which no major disturbances to the plasma from such mechanisms can be expected. For this reason the heating intensity on each side of the target was kept below this limit.

**Comparison of experiment with 1D simulation**

It is interesting to compare the experimental results presented so far with the predictions of the 1-D code presented in Sect.2.1 in order to see the main differences between a 1-D hydrodynamic model and the actual evolution of the plasma. **Fig.4.3.4** shows a comparison between the electron density profiles obtained from the 1-D modelling and the experimental profiles obtained from the interferograms taken at 3.0 and 4.3 ns. The experimental electron density profiles along the $x$-axis of Fig.4.2.1, obtained at 3.0 ns and 4.3 ns after the plasma formation are compared with the analogous calculated profiles presented in Fig.2.1.3. The dashed-dotted portions of the low density tails of the experimental curves were obtained by extrapolation of the experimental data.

**Fig.4.3.4** Comparison between the calculated electron density profiles and the experimental profiles at 3.0 ns and 4.3 ns after the peak of the heating pulse respectively. The dashed and the dashed-dotted portions of the experimental curves are discussed in the text.
The calculated density decreases very slowly with the distance from the target plane, the scalelength ranging smoothly from several millimetres at one target diameter (400 µm) from the target plane, down to several hundred microns at ten times this distance. In contrast, the experimental profile is characterized by three distinct regimes. For distances from the target plane smaller than the target radius, i.e. smaller than 200 µm, the dashed portion of the experimental profile is close to the calculated one although a minor discrepancy in the peak density is evident and will be discussed below. For distances close to the target diameter the electron density decreases very rapidly, the scalelength being of the order of hundred microns or less. Finally, for distances much greater than the target diameter a new regime is approached, indicated by the dashed-dotted portion of the curve, where the density scalelength is similar to that given by the 1-D modelling, i.e. of the order of a few millimetres. However, in this region, the measured density is much lower than the calculated one suggesting that at this distance from the target plane 2-D effects strongly affect plasma expansion.

It is rather surprising however, to find the maximum density higher than the one predicted by the 1-D simulation. The actual laser energy delivered on target was evaluated very carefully and an overestimate of the laser intensity on the target can be excluded. The fact that the electron density is overestimated in the region close to the original target plane can be attributed to a departure of the plasma symmetry from the cylindrical symmetry assumed for the Abel inversion (see Sect.4.2). In fact, since the spot of the heating laser beams is larger than the Al dot, a plasma will be produced from the plastic substrate which is mostly at the top and bottom of the dot (see also the experimental set-up of Fig.4.1.1). For small distances from the target plane, these two plasma slabs can confine the Aluminium plasma in a elongated rather than a pure cylindrical symmetry. This confinement has two consequences. Firstly, it keeps the Al plasma denser than in the case of a free expansion and secondly, it stretches out the plasma in the line of view of the probe beam. This latter fact would give a phase shift bigger than that expected in a cylindrical symmetry, thus leading to an overestimate of the electron density in the region close to the target plane. It is therefore likely to assume that the actual maximum electron density would be not far from the maximum density calculated in a 1-D approximation. This assumption becomes more accurate at earlier times, when the hydrodynamic expansion is more accurately described by a 1D model, and the confining effect of the plasma produced from the plastic is much less significant.

Fig.4.3.5 shows the temporal evolution of the maximum electron density obtained from the 1-D numerical simulation presented in Sec.2.1. According to this
curve the electron density goes below the critical density approximately 200 ps after the peak of the heating pulse. Moreover, at the time at which the delayed interaction typically takes place, the maximum electron density is just below $0.1 n_c @1\mu m$. From an experimental viewpoint, the preformed plasma created with this experimental configuration shows a denser, well localised bulk, and a lower density tail which, 3.0 ns after plasma formation, already extends for a few millimetres. These two regions are separated by a rather steep density gradient. Both the gradient and the bulk play an important role in interaction experiments performed in these conditions and will be discussed in Sect.4.6.

![Temporal evolution of the maximum electron density as obtained from the numerical simulation presented in Sect.2.1. A 250 nm thick Al target was irradiated at a wavelength of 1.053 $\mu m$ and an intensity of $3\times10^{13}$ W/cm$^2$.](image)

An important consequence of the particular longitudinal plasma density profile is that the delayed interaction beam is prevented from propagating freely through the whole plasma. A simple calculation shows that 1 $\mu m$ radiation, propagating in a plasma with a density profile as that shown in Fig.4.3.1 at 3.0 ns, undergoes very weak absorption in the low-density long-scalelength region, but more than 50% of the beam energy is absorbed before the plasma reaches the maximum of the density profile. This is a very profitable condition for studying interaction phenomena in an underdense long scalelength plasma, expanding toward the laser beam. In this case the beam would be almost completely depleted before reaching the plasma located beyond the maximum density, where the expansion velocity is parallel to the $k$-vector of the incident beam.
4.4 - Time-Resolved X-ray Imaging

The study of plasma self-emission in the X-ray region, at photon energies from 100 eV up to several tens of keV is of fundamental relevance to the characterization of laser-produced plasmas. As discussed in detail in Sect.2.2 and Sect.3.4, atomic physics processes are strictly related to the thermodynamic properties of the plasma. Detailed X-ray spectroscopy can therefore provide information on the electron temperature and other relevant plasma parameters. X-ray imaging also allows an effective monitoring of the modifications induced on plasma conditions by laser interaction processes. A description of the techniques used to study X-ray emission from plasmas will be given in this section with emphasis on time-resolved imaging.

X-ray imaging using a pin-hole camera (PHC)

A very useful and easy to use X-ray diagnostic in laser-plasma interaction experiments is the so called pin-hole camera (PHC). A time integrated image of the X-ray emitting region can be generated through a pin-hole as shown in Fig.4.4.1, and recorded on an X-ray sensitive device.

![Fig.4.4.1](image)

Geometry of a pin-hole camera for X-ray imaging of plasmas. The pin-hole diameter determines the spatial resolution of the image. In the limit of high magnification the spatial resolution of the image in the object plane is simply given by the pin-hole diameter itself.
X-ray sensitive films can be used as detectors, however, higher sensitivity and better flexibility from an operative viewpoint can be achieved using an intensified phosphor screen instead. Presently available intensifiers based on the microchannel-plate technology (see below in this section) can easily provide a thousand-fold overall gain in intensity. A commercial B&W film is then used to record the intensified image. According to the geometry of Fig.4.4.1 the magnification of a pin-hole camera is given by \( M = A' / A = q/p \). The spatial resolution of the image is generally given by the combination of geometrical and diffraction effects. It is easy to verify that the size of the image of a point source in the geometrical limit is given by \( \Delta x_g = d_{ph} (1 + M) \), \( d_{ph} \) being the pin-hole diameter. According to the Fraunhofer diffraction theory (Born & Wolf, 1970), the typical transverse size of the Airy pattern produced by a uniformly illuminated circular aperture of diameter \( d_{ph} \) at a distance \( q \), i.e. the diameter of the circle containing 85% of the total energy, is \( \Delta x_d = 2.44 \frac{q \lambda}{d_{ph}} \) , \( \lambda \) being the X-ray wavelength. The best pin-hole size is given by the condition \( \Delta x_g \approx \Delta x_d \), that is,

\[
d_{ph} = \frac{2.44 \lambda q}{1 + M}.
\] (4.4.1)

For typical conditions of laser-plasma experiments, \( \lambda \approx 5 \text{Å} \) , and \( q \approx 50 \text{cm} \), the best pin-hole diameter is approximately 7 \( \mu \text{m} \) for \( M \approx 10 \) and 5 \( \mu \text{m} \) for \( M \approx 30 \). On the other hand a lower limit to the pin-hole size that can be employed in laser-plasma experiments is also set by the available X-ray flux as well as by the manufacturing feasibility. Typical commercially available pin-holes have diameters as small as 5 \( \mu \text{m} \). Therefore, for a magnification \( M < 30 \) the image resolution is determined by the geometrical limit. The size of the smallest feature which can be resolved in the object plane, i.e. in the plasma, is then

\[
\Delta x_{\text{plasma}} \equiv d_{ph} (1 + \frac{1}{M}). \tag{4.4.2}
\]

In the limit of high magnification the spatial resolution of the pin-hole camera image in the object plane is given by the pin-hole diameter. The pin-hole diameter also determines the X-ray flux collected by the camera. The number of 1KeV photons reaching the recording device per unit area of the magnified image can be written as follows

\[
\Phi_{1\text{keV}} = 6.25 \times 10^{15} \frac{\eta_s E_L d_{ph}^2}{16 \pi q^2 M^2 A^2} \text{photons / cm}^2 \tag{4.4.3}
\]
$E_L$ being the laser energy absorbed by the plasma in Joules and $\eta_x$ being the X-ray conversion efficiency. In the conditions of our experiments the absorbed laser energies $E_L$ range from several joules to hundred joules and the X-ray conversion efficiency in the 1KeV region is of the order of $\eta_x \approx 0.1$. With a plasma transverse size $A \approx 100 \ \mu m$ and assuming an image magnification $M = 10$, with a 10 $\mu m$ pin-hole diameter and a pin-hole to image distance, $q \approx 50 cm$, we obtain an X-ray flux ranging from $10^5$ to $10^8$ photons / cm$^2$. This value should be compared with the flux required to obtain an optical density above fog of $D = 1$ on a typical X-ray film (Rocklett et al., 1985), that is, approximately $\Phi_{1keV} = 10^8$ photons / cm$^2$. Therefore the use of an active (intensified) pin-hole camera becomes necessary in these experiments in order to cover the whole range of experimental conditions.

**Time-resolved x-ray imaging**

A crucial aspect of laser-plasma interaction experiments is that the event under investigation develops in a very short time, typically less than a nanosecond. A pin-hole camera would only provide an X-ray image of the plasma integrated over the whole event. In particular, in the case of laser interactions with preformed plasma described in the following sections it is vital, for example, to be able to distinguish between the phase of plasma formation from the actual interaction which usually takes place a few nanoseconds later. Using an X-ray streak-camera one can take a 1D portion of the image and resolve it in time. This technique has been successfully employed for several years.

However it is now possible (Kilkenny, 1991) to temporally resolve fully 2D images with a temporal resolution as small as 50 ps using microchannel-plates (Wiza, 1979) (MCP) as X-ray sensitive devices. These plates are characterized by a very low capacitance that enables fast switching of the driving electric fields. These so called proximity focused X-ray devices can be used as a gated intensifier as well as a DC intensifier and usually do not require additional intensifiers due to their intrinsic high gain. Special configurations can also be designed in order to obtain a sequence of images on the same MCP.

One of these devices, in the following referred to as Strip Line Imager for X-rays (SLIX), has been tested and then successfully employed in our experimental investigation on long-scalelength plasmas. It was designed to allow four-frame X-ray imaging with 120 ps gate-time with the inter-frame time adjustable from zero to a few nanoseconds. A diagram of the SLIX sensitive unit is shown in Fig.4.4.2. The input surface of the microchannel-plate detector is coated with a 500nm copper
layer in four separated rectangular regions (Strip-Lines). These strip-lines can be independently driven when a voltage is applied between them and the MCP output surface, which is uniformly coated with a 500 nm thick copper layer. A photon incident on the MCP input surface generates photoelectrons that, accelerated by the external electric field, hit the walls of the micro–channel producing secondary electron emission. This process takes place several times during the flight of the electrons in the micro–channel resulting in a photomultiplier-like gain.

![Diagram of Strip-Line X-ray Imager detection system](image)

**Fig. 4.4.2** Schematic arrangement of the Strip-Line X-ray Imager detection system with the MCP input surface showing the four 500 nm thick Cu coated regions.

The electrons produced are finally driven onto a phosphor screen by a second electric field. Due to the small electron path distances (see Fig. 4.4.2), a relatively large electric field is capable of preserving the spatial distribution of the outcoming electrons. A standard B&W film records the light produced by the phosphor screen and collected by a fibre optic plate.

The SLIX was coupled to a specifically designed pin-hole camera able to produce four identical images of the plasma onto the four sensitive strip-lines of the MCP. The PHC magnification and spatial resolution were chosen in order to resolve details of about 10 µm in the plasma. Since the intrinsic SLIX spatial resolution, determined by the MCP structure, is about 100 µm, a minimum image magnification of 10 was required to resolve 10 µm structures in the plasma. In order to produce four identical images onto the MCP input surface as shown in **Fig. 4.4.3**, a four pin-hole array was used, whose geometry was set by the strip-line position on the MCP and the PHC image magnification. Because of the very small distances (fractions of mm) involved, the single pin-hole mount shown in **Fig. 4.4.3**
was made by drilling the four pin-holes on a 30 µm thick Platinum foil using a Q-switched 40 ns Nd laser beam focused onto the foil via a microscope objective. The PHC magnification in the final set up used for the experiments described below was approximately 13 with a pin-hole size close to 10 µm. The X-ray filter used in the typical configuration was a 3 µm thick Al foil.

![Fig.4.4.3. Pin-hole arrangement for the four-frame gated X-ray Imager SLIX. The geometry of the pin-hole array is defined once the image magnification is chosen. With 13x magnification the four pin-holes lay on the vertices of a 700µm square.](image)

Since this device was usually employed to detect X-ray emission from Al plasmas, it is important to examine the behaviour of the filter relative to the spectral distribution of the radiation under investigation. Fig.4.4.4 shows the transmittivity of the 3 µm Al filter in the spectral range of interest in this study.

Also plotted on the same graph is the X-ray spectral emissivity calculated using the atomic physics code RATION (Lee et al., 1984), in collisional-radiative equilibrium, for an Al plasma at an electron density of $1.2 \times 10^{20}$ cm$^{-3}$ and an electron temperature of 400 eV. Due to the K-shell absorption edge at 1560 eV, most of the line emission from He-like and H-like ions is strongly absorbed. The dominant contribution to the transmitted X-rays will therefore come from continuum radiation, namely free-bound and free-free. These circumstances are of particular interest from a diagnostic viewpoint. In fact continuum radiation, and in particular that originating from free-free processes, is closely related to the electrons and is only coupled to the ions via their charge number and the electron-ion collision frequency. In contrast with bound-bound radiation which is related to internal atomic processes as well as to the particular equilibrium holding in the plasma among ions, electrons and radiation, free-free radiation is relatively independent from these factors. It can therefore offer a more direct approach for investigating radiation induced modifications of plasma hydrodynamic conditions, such as those arising from local energy deposition, whereas these modifications occur on a time-
scale short compared to the typical time-scale of relaxation of atomic physics processes, as discussed in Sect.3.4.

![X-ray transmittivity of the 3 µm Al filter used in the SLIX imaging system.](image)

**Fig.4.4.4.** X-ray transmittivity of the 3 µm Al filter used in the SLIX imaging system. The X-ray emissivity of a 400 eV, $1.2 \times 10^{20}$ cm$^{-3}$ Al plasma in collisional-radiative equilibrium is also plotted for comparison.

X-ray images of the plasma were taken irradiating thin Al targets in the line-focus geometry described in Sect.4.1. The experimental set-up is schematically shown in Fig.4.1.2. **Fig.4.4.5** shows a time integrated X-ray image of an Al target heated with 600 ps, 0.53 µm laser pulses at an intensity of $\approx 2 \times 10^{13}$ W/cm$^2$ on each side, obtained with the SLIX working as a DC X-ray intensifier. The most striking feature of this image is that the structure of the X-ray source, even though integrated in time, clearly matches the shape of the heated target, as shown in Fig.4.1.1, consisting of a 500 µm long, 300 µm wide, 700 nm thick Al film supported by a 100 nm thick Formvar substrate. The X-rays emitted by the heated plastic substrate are also visible at the edges of the main Al stripe. The feature on the left-hand side of the image, elongated along the vertical direction, is the X-ray emission produced by the delayed interaction pulse which reaches the plasma 2.2 ns after the heating pulses.

The broad-band 1.053 µm, 600 ps laser pulse used for interaction purposes was smoothed using the induced spatial incoherence (ISI) technique which resulted in a slightly longer pulse, $\approx 800$ ps. The effect of ISI on the intensity distribution in the focal plane as well as its consequences on interaction processes will be discussed in detail in Sect.4.6. The beam was focused into the plasma in a 140 µm (FWHM) focal spot using an f/10 optics and the intensity was approximately $10^{13}$ W/cm$^2$. 

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The X-ray emission corresponding to the interaction process is confined in a small region, approximately 300 µm wide and ≈50 µm long, and is located approximately 50 µm from the edge of the Al stripe. This fact clearly suggests that lateral plasma expansion, i.e. expansion parallel to the target plane has occurred. This observation supports the conclusion formulated in Sec.4.3 concerning two-dimensional effects, which were found to play an important role in the hydrodynamic expansion of the plasma.

![Image](image_url)

**Fig.4.4.5** Time integrated X-ray image of a 700 nm thick, 500 µm long, 300 µm wide Al target heated on each side by a 600 ps, 0.53 µm laser pulses at an intensity of $\approx 2 \times 10^{13}$ W/cm$^2$. The image was obtained with the SLIX working as a DC X-ray intensifier.

The effect of the delayed interaction on the plasma conditions is clearly seen in time-resolved images where the various phases of the event can be studied separately. Temporal resolution can be achieved on the SLIX device when a gated voltage is applied to each strip-line. A 650 V, 200 ps FWHM pulse was used to generate an electric field of approximately 15 KV/cm in the MCP, which was added to a DC bias voltage ($\leq 300$ V). Due to the extremely non-linear behaviour of the MCP output current with the applied electric field, the gain produced by the low DC electric field is negligible. Nevertheless, once added to the 200 ps gating pulse, this DC field activates the gain of the MCP. On the other hand by varying the DC voltage it is possible to adjust the overall MCP amplification. The gate pulse generated by a single pulse generator was split into four identical pulses which were applied to the four strip-lines using high frequency cables whose length was varied in order to synchronize the four frames relative to each other.

Images like the one shown in Fig.4.4.5 were produced on each of the four frames of the SLIX and then, by driving the frames at different times, a temporal sequence of X-ray images of the laser-plasma interaction region was recorded. The timing of each frame was set according to the particular configuration of the interaction.
experiment as shown schematically in Fig.4.4.6. Typically the first frame was synchronized with the heating laser pulses (H) while the remaining three frames were usually synchronized on the interaction laser pulse (I). The delay between the heating and the interaction was typically 2.2 ns.

A sequence of time resolved X-ray images of the laser-plasma interaction region is shown in Fig.4.4.7. A target similar to that of Fig.4.4.5, i.e. 500 µm long, 300 µm wide, 700 nm thick Al film supported by a 100 nm thick Formvar substrate, was heated at an intensity of $\approx 4 \times 10^{13}$ W/cm$^2$ on each side. The intensity of the 600 ps interaction pulse was $5.8 \times 10^{13}$ W/cm$^2$ and its delay relative to the peak of the heating pulse was 2.2 ns. As indicated in Fig.4.4.6, frame #1 was synchronized with the peak of the 600 ps heating pulse while the remaining three frames (#2, 3 & 4) were timed to record the evolution of the interaction process with the preformed plasma. According to the numerical analysis presented in Sect.2.1, the maximum electron density, at the peak of the heating pulse (See Fig.2.1.4) is expected to be slightly above the critical density for the 0.5 µm heating laser light ($n_{c@0.5\mu m} = 3.8 \times 10^{21}$ cm$^{-3}$). At the time of the delayed interaction, i.e. 2.2 ns later, the maximum electron density is expected to be $1.5 \times 10^{20}$ cm$^{-3}$. On the other hand, from the comparison of experimentally measured electron density profiles and calculated ones, as already pointed out in Sec.4.4, the maximum density is expected to be reasonably well predicted by the 1 D hydro-code used in Sect.2.1. The X-ray emission relative to the heating phase, though very strong in the first frame, goes below the detection level of the X-ray imaging system in the following frames. This is consistent with the X-ray time resolved spectra taken in similar experimental conditions and presented in Sect.4.5. In that section a simple model of X-ray emissivity is presented and discussed to explain the temporal behaviour of experimental X-ray emission (See Fig.4.5.6).
4.4 Time-Resolved X-ray Imaging

Intense X-ray emission takes place when the interaction pulse strikes the preformed plasma. This emission is localized in a rather small region whose transverse size is approximately 200 µm, i.e. comparable to the focal spot size of the interaction beam (140 µm FWHM). Also, taking into account the angle of view (45 degrees) of the SLIX device with respect to the target plane as shown in Fig.4.1.2, the actual longitudinal size of this region varies from 300 µm in the frame #2 to approximately 360 µm in the frame #3. The emissivity of this region strongly decreases with time, at the end of the interaction pulse as shown by the weak image of frame #4. These circumstances, which make the interaction region “visible” in the X-rays, indicate that a strong absorption of laser energy by the plasma takes place in this small region. Since the maximum plasma electron density is well below the critical density for the infrared interaction beam, \( n_e/n_c \approx 0.14 \), the absorption can be mainly attributed to collisional absorption.

**Inverse bremsstrahlung absorption in preformed plasmas**

The longitudinal extent of the X-ray emitting region is therefore expected to be related to the IB absorption length. According to Eq.1.2.13 and assuming propagation of the laser e.m.wave in a uniform plasma with temperature and density given by the simulation shown in Fig.2.1.5, i.e. \( T_e = 450\text{eV} \) and \( n_e = 1.5 \times 10^{20} \text{ cm}^{-3} \) the absorption length for IB absorption is predicted to be

![Fig.4.4.7 Sequence of time resolved X-ray images of laser interaction with a preformed plasma produced by irradiation of a 700 nm thick, 300 µm wide, 500 µm long Al stripe at an intensity of \( \approx 4 \times 10^{13} \text{ W/cm}^2 \) on each side. The interaction beam was focused on the left edge of the plasma in a 140 µm (FWHM) focal spot at an intensity of \( 5.8 \times 10^{13} \text{ W/cm}^2 \). The timing of each frame relative to the laser pulses is schematically shown in Fig.4.4.6.](image-url)
approximately 500 µm which is approximately a factor 1.4 greater than the length of
the X-ray emission region. A more direct comparison between these two lengths
would require the dynamics of X-ray emission processes to be accounted for as a
function of many effects including electron heat transport, spatial beam energy
depletion, non-linear laser-induced instabilities like stimulated Raman scattering and
two plasmon decay. In addition we notice that, due to the strong dependence of the
IB absorption coefficient on the electron density (see Eq.1.2.13), $\kappa_{ib} \propto n_e^2$.
Therefore, a 25% error in the estimated value of the electron density would make the
calculated absorption length consistent with the measured value of 360 µm.

A strong dependence of the interaction length upon electron density was also
observed experimentally. Preformed plasmas with lower electron densities were
generated using thinner Al targets, i.e. 500 nm thick instead of the 700 nm of
Fig.4.4.7. In this case, according to the simulation, the maximum electron density is
expected to be approximately $1.1 \times 10^{20}$ cm$^{-3}$. The absorption length at the
predicted electron temperature of 400 eV is calculated to be 1.2 mm, which is larger
than the length of the preformed plasma itself. In fact, this prediction was
experimentally confirmed as the interaction length in this case, matched the length of
the plasma, showing no local heating effects as in Fig.4.4.7.

**Heat transport and X-ray emission**

While the longitudinal extent of the X-ray emitting region gives insight into the
absorption processes, the transverse extent and in particular its dependence upon the
intensity of the interaction beam, can be related to electron heat transport effects.
**Fig.4.4.8** shows a sequence of time resolved X-ray images of the laser-plasma
interaction region obtained in similar conditions to that of Fig.4.4.7 but at a lower
interaction intensity of $3.5 \times 10^{13}$ W/cm$^2$. The transverse size of the interaction
region as shown in frame #3 is now approximately 150 µm compared to the 200 µm
of frame #3 in Fig.4.4.7. It should be noted however, that the value of the intensity
of the interaction beam on the plasma is only an average value calculated assuming
that the beam energy is uniformly distributed in the 140 µm (FWHM) focal spot
during the 600 ps pulse duration. The effect of the energy in the “wings” of the
focal spot is therefore not accounted for by this simple description. At increasing
energies these wings will start to contribute to the transverse size of the interaction
region as the local laser intensity, although a fraction of the intensity in the FWHM,
is high enough to heat the plasma and give rise to X-ray emission.
A better understanding of this effect can be achieved by studying its dependence upon the nominal laser intensity. The plot of Fig.4.4.9 shows the transverse size of the interaction region as measured from time resolved X-ray images as a function of the interaction intensity in a range of from \(10^{13}\) to \(\approx 10^{15}\) W/cm\(^2\) obtained with the same focusing conditions, i.e. 140 µm FWHM focal spot. The two sets of data were obtained with two different preformed plasma. Al targets as well as KCl targets were used whose average atomic numbers are 13 and 18 respectively and average atomic mass are 27 and 37 respectively. A linear fit to the two sets of data is also plotted; notice that a log scale has been used for the intensity axis.

The first remarkable feature of this graph is that, as the intensity decreases, the measured transverse size approaches, in both cases, a value corresponding to the FWHM of the focal spot. In contrast, as the intensity increases, the transverse size increases approximately linearly in both cases as indicated by the two curves plotted, although with two different slopes. However we observe here that the correspondence between the low intensity limit of the measured transverse size and the nominal FWHM is likely to be accidental. In fact in order to explain this behaviour we must take into account the actual intensity profile of the focal spot. For simplicity we will consider a gaussian intensity distribution in the far field of the interaction beam. For a given average power in the pulse there will be a corresponding intensity distribution in the focal spot with a maximum at the centre of the spot. We can reasonably assume that the local heating effect will give rise to detectable X-ray emission in the selected spectral range only if the local power is
higher than a given reference value. Therefore an effective size of the gaussian focal spot can be calculated which, in general, is different from the FWHM. As the average power of the pulse increases, the effective transverse size of the focal spot increases.

![Graph showing the transverse size of the X-ray emitting region as a function of nominal intensity.](image)

**Fig. 4.4.9** Transverse size of the X-ray emitting region as a function of the nominal intensity of the interaction beam focused on an Al plasma and a KCl plasma. The plasma heating intensity was approximately $4 \times 10^{13} \text{ W/cm}^2$ on each side of the target. The solid curve shows the behaviour of an effective spot size calculated for a gaussian beam (see text).

The dependence of this effective size upon the nominal intensity has been calculated assuming a gaussian intensity profile and a reference intensity of $10^{12} \text{ W/cm}^2$ which was found to be the typical value of the nominal intensity below which no X-ray emission was detected. According to this result shown by the solid line in Fig. 4.4.9, as the nominal intensity decreases, the effective transverse size of the focal spot decreases and vanishes when the maximum power approaches the reference value. This is in contrast with the behaviour of the measured size of the X-ray emitting region, which, for nominal intensities below $10^{14} \text{ W/cm}^2$, saturates to approximately 150 µm.

In order to compare the calculated size with this measured one we must consider the effect of heat transport phenomena which give rise to a spread of the thermal energy over a volume determined by the characteristic thermal diffusion length. When the calculated effective size is smaller than this length, thermal effects will play a dominant role in determining the size of the X-ray emitting region. Using a 1D equation of classical thermal transport one can find (Afshar-rad et al., 1992) the characteristic thermal diffusion length, $L_D \equiv (\kappa_{SH} t / n_e)^{1/2}$ as a function of the time $t$, where $\kappa_{SH}$ is the Spitzer-Härm conductivity given in Sect. 2.4. Assuming a fully
ionized Al plasma at an electron temperature and density of 450 eV and $1.5 \times 10^{20}$ cm$^{-3}$ respectively, and integrating over the 200 ps gate time of the X-ray imaging system, one finds $L_D \approx 175$ µm which is, in fact, very similar to the lower limit of the measured transverse size given in Fig.4.4.9.

The observed different behaviour of the Al and the KCl data can be explained taking into account ionization effects. In fact, as will be shown in detail in Sect.4.5, at the given electron temperature of 450 eV the Al plasma can be considered almost fully ionized. Consequently all the energy transferred from the interaction beam to the plasma is converted into thermal electron energy. On the contrary, in the case of KCl, due to the higher atomic number of the Chlorine, the plasma is far from being fully ionized. In this case further ionization will occur during the interaction process which will subtract thermal energy from the electrons. Consequently, in the case of the KCl plasma the transverse size of the X-ray emitting region is expected to be smaller than in the case of the Al plasma. In the limit of high intensity or, alternatively, high temperature, both plasmas can be considered fully ionized and are expected to give rise to an X-ray emitting region characterized by the same transverse size as given by the solid curve of Fig.4.4.9.

The experimental results presented so far in this section show that time-resolved X-ray imaging of the laser plasma interaction region can provide valuable information of the plasma condition which can be directly compared with simple analytical models or with the predictions of numerical simulations. In addition, the physics of the interaction can also be investigated. The absorption length inferred from the X-ray images presented here was found to be in agreement with the classical theory of collisional absorption. Furthermore, the thermal diffusion length estimated from the data was found to be consistent with the classical thermal transport theory. As will be clear in the analysis presented in Sect.4.6, these results are also consistent with the conclusions of Sect.2.4 which ensure that, in the interaction conditions investigated here, non-local effects in the heat transport process are expected to give little contribution. However, a better understanding of these processes requires a more detailed knowledge of plasma conditions. Time resolved X-ray spectroscopy will be considered in the next section in order to obtain information on the plasma electron temperature.
4.5 - History of Electron Temperature

X-ray line emission from H-like lines and He-like ions has been investigated in order to measure the temporal evolution of the electron temperature of the Al plasmas studied in the previous sections. Line ratios between resonance line emission from He-like and H-like Al ions was taken into account and compared with the calculated ones for a given temperature and density in collisional-radiative equilibrium. According to this model, line intensities depend upon electron density and temperature. Consequently, the measurements of electron density described in Sect.4.3 play a crucial role in the determination of the electron temperature. Opacity effects have also been accounted for by comparing experimental data with the prediction of numerical simulations.

Time resolved X-ray spectroscopy

Time resolved spectroscopy has been employed extensively in our experiments in order to accurately characterize plasma conditions (Gizzi et al., 1994). Experimental data taken in the cylindrical geometry configuration (see Fig.4.1.2) are presented and discussed in this section. As far as the plasma production and characterization in terms of electron temperature is concerned, the cylindrical geometry and the line focus geometry described in Sect.4.1 can be considered very similar to each other. In fact the typical target and laser parameters differ only slightly as shown in Table 4.1.1. Therefore the results obtained for the cylindrical geometry can be easily extended to the line focus geometry.

The temporal evolution of the electron temperature has been inferred by comparison of the experimental line ratios with the predictions of the atomic physics code RATION (Lee et al., 1984) in steady-state collisional-radiative equilibrium as discussed in Sect.2.2. Since the plasma was generated by exploding Aluminium targets, line emission from highly stripped Al ions was considered. In particular, in the plasma conditions of our experiments, He-like and H-like species were dominant. The wavelength and the corresponding photon energy of the most intense resonance lines from these ionic species are reported in Table 4.5.1.
### Table 4.5.1

<table>
<thead>
<tr>
<th>He</th>
<th>$\lambda$(Å)</th>
<th>$\nu$(eV)</th>
<th>TRANS.</th>
<th>H</th>
<th>$\lambda$(Å)</th>
<th>$\nu$(eV)</th>
<th>TRANS.</th>
</tr>
</thead>
<tbody>
<tr>
<td>He$\alpha$</td>
<td>7.75</td>
<td>1600</td>
<td>1s$^2$ - 1s2p</td>
<td>Ly$\alpha$</td>
<td>7.17</td>
<td>1729</td>
<td>1s - 2p</td>
</tr>
<tr>
<td>He$\beta$</td>
<td>6.64</td>
<td>1867</td>
<td>1s$^2$ - 1s3p</td>
<td>Ly$\beta$</td>
<td>6.05</td>
<td>2050</td>
<td>1s - 3p</td>
</tr>
<tr>
<td>He$\gamma$</td>
<td>6.31</td>
<td>1965</td>
<td>1s$^2$ - 1s4p</td>
<td>Ly$\gamma$</td>
<td>5.74</td>
<td>2160</td>
<td>1s - 4p</td>
</tr>
<tr>
<td>He$\delta$</td>
<td>6.17</td>
<td>2010</td>
<td>1s$^2$ - 1s5p</td>
<td>Ly$\delta$</td>
<td>5.60</td>
<td>2214</td>
<td>1s - 5p</td>
</tr>
<tr>
<td>He$\epsilon$</td>
<td>6.10</td>
<td>2032</td>
<td>1s$^2$ - 1s6p</td>
<td>Ly$\epsilon$</td>
<td>5.53</td>
<td>2242</td>
<td>1s - 6p</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>H</th>
<th>$\lambda$(Å)</th>
<th>$\nu$(eV)</th>
<th>TRANS.</th>
</tr>
</thead>
<tbody>
<tr>
<td>Ly$\alpha$</td>
<td>7.17</td>
<td>1729</td>
<td>1s - 2p</td>
</tr>
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<td>Ly$\beta$</td>
<td>6.05</td>
<td>2050</td>
<td>1s - 3p</td>
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<td>Ly$\gamma$</td>
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<td>Ly$\delta$</td>
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<td>1s - 5p</td>
</tr>
<tr>
<td>Ly$\epsilon$</td>
<td>5.53</td>
<td>2242</td>
<td>1s - 6p</td>
</tr>
</tbody>
</table>

Wavelength and corresponding photon energy of resonance lines from He-like and H-like Aluminium. The resonance series of H-like Al ions is labelled according to the analogous Lyman series of the Hydrogen atom.

An X-ray spectrometer consisting of a flat TIAP (TlIHC$_4$H$_4$O$_4$) crystal ($2d = 25.9\text{Å}$ ) was set in a first order Bragg configuration as shown in Fig.4.5.1. An intensified X-ray streak-camera fitted with a CsI photo-cathode was coupled to the spectrometer. A spectral region from approximately 5.6 Å to 6.8 Å was selected in order to study resonance line emission from He$\beta$ to Ly$\delta$.

![TIAP Crystal Configuration](image)

**Fig.4.5.1**. Schematic of flat TIAP crystal configuration in the X-ray spectrometer for time-resolved spectroscopy measurements.

The temporal resolution was determined by the width of the fixed slit in front of the streak-camera cathode and was, depending on the selected streak speed, 50 ps or 100 ps as specified below. The spectral resolution, basically due to the size of the emitting region, was evaluated directly from calibrated X-ray spectra and was found to be typically 40 mÅ. This value is in agreement with a simple calculation made assuming a plasma emitting region whose size is that given by the cross-section of the Al dot target along the line of sight of the spectrometer.
Experimental results

Fig.4.5.2 shows a time-resolved spectrum obtained in the same event as the interferogram of Fig.4.2.2, at a total heating irradiance of $6.5 \times 10^{13}$ W/cm$^2$. The optical density distribution on film was converted into intensity of the incident radiation on the film using the calibration step wedge.

Fig.4.5.2. Time-resolved X-ray spectrum of K-shell Aluminium emission from plasma heated at an irradiance of $6 \times 10^{13}$ W/cm$^2$ by four 600ps FWHM, 1µm laser pulses.

Fig.4.5.3 shows a 1D trace taken 500 ps after the peak of X-ray emission and integrated over 50 ps, which is the temporal resolution of the spectrum. Emission lines from the He$\beta$ to the Ly$\delta$ are clearly visible with the Ly$\gamma$ and Ly$\delta$ emerging from the He-like continuum. The He$\epsilon$ line is just visible as a shoulder-like feature on the long wavelength side of the Ly$\beta$ line.

Fig.4.5.3. Lineout of the spectrum of Fig.4.5.2 taken 500 ps after the peak of X-ray emission intensity and integrated over 50 ps, i.e. over the temporal resolution of the spectrum.
Line intensity profiles were found to be well fitted using gaussian profiles with $\Delta \lambda_{\text{FWHM}} = 40\text{mÅ} \pm 10\%$. Except for the He$\alpha$ and Ly$\beta$ lines all remaining lines are well resolved and it is therefore relatively simple to determine the line intensity. With the available spectral resolution, the Ly$\beta$ line is only partially resolved, being merged with the He$\alpha$ and higher quantum number He-like lines and with the He-like continuum edge. Intensity ratios involving the Ly$\beta$ transition will consequently be affected by a larger error. A better condition from this point of view is accomplished at higher heating laser intensities as hotter plasmas are generated with more intense H-like lines. This effect is clearly shown in Fig.4.5.4 which shows a time-resolved spectrum obtained in similar conditions as Fig.4.5.2 but with a heating laser intensity of $1.2 \times 10^{14} \text{ W/cm}^2$.

![Fig.4.5.4](image_url)  
*Fig.4.5.4. Time-resolved X-ray spectrum of K-shell Al emission from plasma heated at an irradiance of $1.2 \times 10^{14} \text{ W/cm}^2$ by four 600 ps FWHM, 1µm laser pulses. The effect of the interaction pulse delayed by 2.5 ns with respect to the peak of the heating pulse is also visible.*

Also visible in this spectrum is the X-ray emission arising from the delayed interaction pulse set to reach the plasma 2.5 ns after the peak of heating pulses. The effect of the delayed interaction on the conditions of the preformed plasma will be discussed in detail later in this section. However, it is important to notice that the spectral resolution for this part of the spectrum is clearly improved due to the smaller spatial extent of the X-ray source. In fact, in this case, the X-ray emission comes from the region of plasma heated by the interaction beam whereas the initial part of the spectrum comes from the whole plasma extent. This explanation is supported by the time-resolved images presented in Fig.4.4.7 and Fig.4.4.8 where the features of the X-ray emitting region during the plasma heating phase and during the delayed interaction are clearly visible.
The effect of the higher heating intensity compared to the spectrum of Fig.4.5.2, and the consequent increase of the electron temperature is shown by the stronger Lyγ and Lyδ lines which now clearly emerge from the He-like continuum. These circumstances make X-ray spectra like this, obtained at higher laser intensities, more suitable for temperature measurements from line intensity ratios between H-like and He-like Al lines. On the other hand, the conclusions on the electron temperature obtained in these conditions can be reliably extended to the regime of lower heating laser intensities by comparing the experimental results with the hydrodynamic numerical simulation. The Lyγ to Heγ intensity ratio of the spectrum of Fig.4.5.4, restricted to the heating phase only, has been measured as a function of time and the result is shown on Fig.4.5.5.

![Figure 4.5.5](image)

**Fig.4.5.5.** Experimental intensity ratio of the H-like Al γ-line (1s - 4p) to the He-like Al γ-line (1s^2 - 1s4p) as a function of time with respect to the peak of the He-like Al β-line.

The zero of the temporal axis corresponds to the time at which the He-β line, that is the most intense line in the spectrum of Fig.4.5.4, reaches its maximum. Incidentally, we notice that the lines shown in the spectrum, were found to reach their peak approximately simultaneously. Only a small difference could be detected between the set of the He-like lines and the set of the H-like lines, the latter peaking a few tens of picoseconds later. This particular choice of origin for the temporal axis will be justified by the analysis presented later in this section. In fact, according to the simulation, the peak of the X-ray emission is very closely related to the peak of the heating laser pulse, which is the main temporal *fiducial*.

As already pointed out, spectra like that shown in Fig.4.5.4 result from spatial integration over the whole plasma extent. Consequently plasma parameters determined by analysis of line emission spectra should, in principle, be regarded as
averaged over plasma electron temperature and density distribution in space. However, since line emission intensity is, by itself, strongly dependent on local electron temperature and density, we can expect that particular plasma regions will give dominant contributions to the measured line emission.

**Effects of spatial integration**

A simple estimate of such an effect can be made in the limit of coronal equilibrium in an optically thin steady-state plasma. This approximation is likely to be unrealistic at the beginning of the heating pulse when density of the emitting plasma is too high and laser intensity increases too rapidly for a thin steady state plasma approximation to be fulfilled. However, later in time, typically after the peak of the laser pulse, the electron temperature is high enough, and the electron density low enough that we can assume a plasma equilibrium not far from coronal equilibrium. In fact, according to the plot of Fig.2.2.1, an hydrogen-like Al plasma at an electron temperature of 500 eV can be considered in coronal equilibrium if the electron density is of the order of \(10^{19}\) cm\(^{-3}\) or less. According to the experimental results summarized in Fig.4.3.4, the electron density of the plasma relative to the spectrum of Fig.4.5.2 ranges from \(10^{18}\) to \(10^{20}\) cm\(^{-3}\). Therefore, most of the plasma volume is expected to be not far from coronal equilibrium. However, the deviations from the coronal limit are expected to affect the intensity ratios and to have only minor consequences on the general features of the X-ray emission including the spatial distribution and the temporal behaviour. We will therefore apply this model to the whole plasma evolution even though the comparison with experimental results, particularly in the early stages of plasma formation will have to be carefully thought about.

Using the general expression for the intensity of a line in an optically thin coronal plasma given by Eq.2.2.12 and substituting the collisional excitation coefficient according to its expression given by Eq.3.3.4, one can find the dependence of the intensity upon the electron temperature and density. If \(n\) is the upper level from which the transition occurs to the ground state \(o\) of a given ion, the intensity of the line depends upon density and temperature according to

\[
I_{n \rightarrow o} \propto n_e n_o T_e^{-1/2} \exp\left(-\frac{\Delta E_{n \rightarrow o}}{T_e}\right),
\]

(4.5.1)

\(\Delta E_{n \rightarrow o}\) being the transition energy, \(n_o\) being the ground state population density of the given ion species. For a given ion charge \(Z\), the intensity of Eq.4.5.1 is a function of hydrodynamic plasma parameters, \(T_e\) and \(n_e\).
The spatially dependent population densities of the ground state of each ionization stage were calculated by the hydro-code MEDUSA using a non-LTE model. According to the simulation, between 90% and 99% of the ions are in their ground state. Therefore only a minor error is made when the average ionization degree is calculated taking into account the population densities of the ions in their ground state. Incidentally, for consistency with the simulation, in Eq.4.5.1 we can reasonably assume that \( n_o \approx n_i \). The intensity of the Heβ line as given by Eq.4.5.1 was obtained post-processing the results of the hydrodynamic simulation. **Fig.4.5.6** shows the emissivity of the Heβ line calculated according to Eq.4.5.1 as a function of distance from the target plane, 250 ps before the peak of the heating pulse. Also shown is the corresponding calculated density profile.

**Fig.4.5.6.** Plasma electron density and Heβ emissivity as a function of distance from target plane as predicted by MEDUSA, 250ps before the peak of the heating pulses.

At this time the plasma is still overdense; emissivity increases going towards higher densities and reaches its maximum in the proximity of the critical density layer. Then it decreases rapidly as electron temperature falls while going towards the supercritical density. The same features can be observed in analogous plots taken later in time. Most of the X-ray emission comes from a narrow plasma region whose temperature and density maximise the emission intensity as given by Eq.4.5.1. Finally **Fig.4.5.7** shows the calculated intensity of the Heβ line integrated over the whole plasma, as a function of time with respect to the peak of heating laser pulse. The result of the simulation is compared with the experimental intensity of the Heβ obtained from the spectrum of Fig.4.5.2. According to this plot, the Heβ line emission reaches its peak approximately 100 ps after the peak of the laser pulse; this is also the time at which the maximum plasma density becomes smaller than the
critical density for the 1 µm heating laser light. The discrepancy between experiment and calculation early during the emission is probably a consequence of the steady-state coronal equilibrium model which, as shown in Sect.2.2 and Sect.3.4, cannot accurately describe this stage of plasma formation.

More significant is the discrepancy observed in the decay of the emission where the experimental curve shows a slower decay rate than the calculated one. This effect has already been observed in a previous experiment (Willi et al., 1989) where time-resolved X-ray spectra from laser plasmas were compared with the results of the simulation performed using a heat flux limiter of 0.1, the same value used for our simulation. It is clear that, later in time, after the explosion of the target, two-dimensional effects, which are not accounted for in our simulation, are expected to play an important role. A more sophisticated modelling of X-ray emission is required here in order to simultaneously deal with both 2D hydrodynamics and atomic physics processes in a self-consistent way. Nevertheless, our calculation, performed in the limit of a steady-state coronal equilibrium and based on 1-D numerical simulations of hydrodynamic plasma expansion, offers a relatively simple way of dealing with the problem and is nevertheless found to be in reasonably good agreement with experimental results.

Further insight into the dynamics of radiation emission is provided by the curves of Fig.4.5.8 which show the temporal behaviour of peak emissivity and the electron density at which such peaks take place. As long as the plasma is overdense, i.e. initially in the laser pulse, the maximum of the emission is located in a plasma region with an electron density ranging from $n_c$ and $3n_c$.
Once the plasma has become underdense, the laser propagates beyond the maximum density layer, heating the whole plasma more uniformly. The electron temperature becomes uniform over the whole plasma extent. Consequently, after this time, dominant contribution to X-ray line emission will come from the peak density region located near the original target plane.

**Opacity effects on temperature measurements**

In order to minimise the contribution of opacity effects, it is always preferable to restrict the analysis to high quantum number members of a resonance series which are less sensitive to opacity effects. On the other hand, line intensity decreases dramatically going towards high quantum number members resulting in low signal to noise ratios. The intensity ratio between the aluminium Ly and He lines considered here represents, in our experimental conditions, the best compromise between these two limits. It will be shown however that, although opacity effects in general play an important role in determining the intensity of X-ray line, at the time of interest for interaction experiments, they can be neglected as they lead to a minor error in the final value of electron temperature.

The intensity ratio between Ly and He was calculated for a uniform plasma as a function of the electron density and temperature and for an optically thin plasma. Moreover, opacity effects were included in the calculation considering lengths of plasma of 10 µm and 100 µm respectively. The results are shown in Fig.4.5.9 for
densities ranging from $10^{18}$ to $10^{22}$ cm$^{-3}$ and for temperatures from 380 to 800 eV. We notice that, for densities below a few times $10^{19}$ cm$^{-3}$ the intensity ratio is almost independent of both density and opacity effects. However, in the range of densities between a few times $10^{19}$ cm$^{-3}$ and a few times $10^{21}$ cm$^{-3}$, there is a strong dependence of line ratios upon density. Also opacity effects strongly affect line ratios only for densities above $10^{20}$ cm$^{-3}$.

![Fig.4.5.9](image.png)

**Fig.4.5.9.** Ly$\gamma$ to He$\gamma$ intensity ratios as a function of electron density for different electron temperatures. Solid line curves have been obtained with no opacity included in the calculation while dashed line and dashed-dotted lines correspond to calculation with opacity effects included and with 10 µm thick and 100 µm thick plasma respectively.

According to these results, electron density and opacity are expected to play a dominant role in determining intensity ratios during the explosion of the target when the density is high. However, later in time, and in particular at the time of interest for laser-plasma interaction physics, i.e. few nanoseconds after the peak of the laser pulse, plasma conditions will be such that opacity effects will affect the Ly$\gamma$ to He$\gamma$ intensity ratio only slightly. We notice that the region of interest for X-ray line emissivity is always located near the target plane, where plasma hydrodynamics is expected to have a 1D behaviour. Therefore, the plot of Fig.4.5.8, obtained in the 1D approximation, can provide the additional information on the density of the emitting region required to obtain the temperature from the experimental line ratio.

**Fig.4.5.10** shows the temporal dependence of the electron temperature obtained comparing the experimental intensity ratios of Fig.4.5.5 with the plots of Fig.4.5.8 and Fig.4.5.9. The origin of the time was set according to the simulation, as in the case of Fig.4.5.7. Each curve of the graph was obtained comparing the experimental ratio with the calculated one and assuming an opacity effect with a given length of
homogeneous plasma and looking for the electron temperature at which agreement between experiment and theory was achieved. The same procedure used for the three plasma lengths taken into account in Fig.4.5.9.

![Graph showing electron temperature over time with different opacity conditions.](image)

**Fig.4.5.10.** Temporal evolution of the electron temperature of the plasma obtained comparing the experimental intensity ratio of Lyγ to Heγ with the prediction of the steady-state Collisional Radiative numerical code RATION.

As already observed, opacity effects give a major contribution during the pulse while later than about 1 ns after the peak of heating pulses, the three curves approach each other and give a well defined electron temperature within less than 100 eV. In particular, at 2 ns, when the electron density is of the order of a tenth of the critical density, the values of the three curves agree within less than 50 eV, giving an average electron temperature of 550 eV, the error being less than 8%.

**Plasma heating by delayed interaction**

The effect of the delayed interaction on the plasma conditions as seen through the time resolved X-ray spectroscopy was introduced when describing the features of the spectrum in Fig.4.5.4. In that spectrum we found that the interaction pulse produced a second burst of X-rays. The peak of this emission was delayed relative to the peak of the first burst, generated by the plasma heating pulses, by approximately the same delay between the heating and the interaction laser pulses.

Moreover, the duration of the X-ray emission was found to be comparable with the 600 ps laser pulse-length. Another feature of the spectrum in Fig.4.5.4 is that there is evidence of a spectral shift, variable in time, of all the lines simultaneously. This is certainly related to the particular configuration of the interaction beam. In fact, in this case the far field of the interaction laser beam consisted (Giulietti et
al., 1993) of a set of three line focal spots, parallel to each other separated by approximately 150 µm, each spot being 100 µm wide and 800 µm long. The average intensity in each of the three spots was approximately $10^{14}$ W/cm$^2$. This particular shape of the focal spot was adopted in order to induce controlled intensity modulations in the interaction region in order to study the effect on the filamentation instability, further discussion of this topic will be given in Sect.4.6. A more direct analysis of the modifications induced by the interaction pulse on the plasma conditions was carried out on the spectrum of Fig.4.5.11 where all the effects described above are more clearly visible due to the higher interaction intensity.

![Time-resolved X-ray spectrum of K-shell Aluminium emission from plasma heated at an irradiance of $1.2 \times 10^{13}$ W/cm$^2$ by four 600 ps FWHM, 1µm laser pulses.](image)

**Fig.4.5.11.** Time-resolved X-ray spectrum of K-shell Aluminium emission from plasma heated at an irradiance of $1.2 \times 10^{13}$ W/cm$^2$ by four 600 ps FWHM, 1µm laser pulses. The interaction beam was focused on the plasma in a 800 µm by 100 µm line focus at an intensity of $3.0 \times 10^{14}$ W/cm$^2$ and was delayed by 2.5 ns with respect to the peak of the heating pulse.

The plasma was produced under similar conditions to that of Fig.4.5.4 and the interaction beam consisted of a single line focal spot as described above with an intensity of $3 \times 10^{14}$ W/cm$^2$. In this case no line shifting was observed and the effect of the improved spectral resolution during the interaction phase, when compared to that heating phase, is much more evident. In fact, the He$\varepsilon$ line can clearly be resolved during the interaction while it is almost completely merged to the Ly$\delta$ line during the heating. **Fig.4.5.12** shows the temporal evolution of the electron temperature obtained from the spectrum of Fig.4.5.11 relative to the interaction phase. In this case, due to the higher intensity available, the temperature was measured taking into account the Ly$\gamma$ to He$\gamma$ intensity ratio as well as the Ly$\delta$ to He$\delta$ intensity ratio. Consistently with the plot of Fig.4.5.10, the error in the data points relative to the late heating phase, that is for $0 < t < 1$ns, arise from opacity...
effects as previously described. In contrast, the large error in the data points relative to the interaction phase, that is for $t > 1 \text{ ns}$, are likely to be related to the spatial inhomogeneity arising from the localized heating of the plasma. In fact, in contrast with the case of Fig.4.5.10, where uniform heating occurred, in this case the plasma has two clearly distinct regions characterized by two different electron temperatures as shown by the time resolved X-ray images of Figs.4.4.7-8 and discussed in detail in Sect.4.4. We can distinguish between the hot plasma region directly heated by the interaction pulse and the cooler background plasma which, as evident in the spectrum of Fig.4.5.11, still contributes to the spectrum and particularly to the He–like components.

**Fig.4.5.12.** Temporal evolution of the electron temperature of the plasma region heated by the interaction pulse obtained comparing the experimental intensity ratio of Fig.4.5.11 with numerical calculations.

Despite the occurrence of these effects, the plot of Fig.4.5.12 clearly shows the change in the electron temperature induced by the delayed interaction pulse. The temperature at the end of the heating phase is consistent with the plot of Fig.4.5.10. In this case, however, the higher interaction intensity starts affecting the plasma already at 1.5 ns, when the electron temperature starts rising again reaching its peak at approximately 2.5 ns. According to the previous analysis, this time is very close to the peak-time of the interaction laser pulse.

This study shows that, in the conditions considered here, the delayed laser pulse interacting with an underdense plasma can give rise to a net increase of the electron temperature in the interaction region of approximately 200 eV. These circumstances should be carefully considered as they can play a very important role in the interaction physics and in particular on the dynamics of laser induced instabilities.
Summary of the main results on characterization

The detailed analysis of the interferometric measurements and the X-ray spectra given so far in this Chapter provides a comprehensive picture of plasmas created by a two-side balanced irradiation of thin Al targets. A full system of interference fringes was visible later than 4 ns after plasma formation, which allowed a complete measurement of electron density at this time. However partially complete patterns were obtained as early as 3 ns and the maximum density could be consistently deduced for earlier times from the numerical simulation. The use of a Fourier transform technique for the analysis of the interferograms allowed the uniformity of the plasma density to be characterized on a scalelength of interest for laser induced instabilities, that is several tens of microns. The comparison of longitudinal experimental density profiles with the ones calculated using 1D numerical simulations shows that, despite the relatively large transverse target size, 2D effects play an important role in plasma expansion. This is also confirmed by the discrepancy, during the final stage of the plasma expansion, between the experimental temporal evolution of X-ray emission and the calculated one.

Time-resolved X-ray spectra, appropriately corrected for opacity, provide the temporal evolution of the electron temperature. The limit to the X-ray spectral resolution was found to be given by the size of the plasma, which, in the case of plasmas for interaction studies is typically rather large, both longitudinally and transversally to the plasma flow. In our experiment, however the line broadening was small enough to allow the electron temperature to be inferred from line intensity ratios. A question arises on where in the plasma most of the X-ray radiation used to infer the electron temperature is generated. It was shown that, at the beginning of plasma formation, X-ray emission comes from a region close to the critical density. However, at the times of interest for interaction studies, when the plasma is well underdense for the laser wavelength of interest, the electron temperature is substantially uniform and X-ray emission comes primarily from the plasma region with the highest density, located near the target plane.

The temporal evolution of the electron temperature was also measured during the interaction phase, when opacity effects on the X-ray line emission were found to be negligible. At 2 ns after the peak of the heating pulses the temperature was found to be 550 eV when the heating intensity on each side of the Al target was of $6 \times 10^{13}$ W/cm$^2$ and was estimated to be 400 eV at an intensity of $3 \times 10^{13}$ W/cm$^2$.

The maximum electron density at this time is expected to be approximately $n_e/10$ in the lower intensity case and $n_e/15$ at the higher intensity.
The interaction of this preformed plasma with a delayed pulse was found to give rise to substantial modifications in the local plasma conditions as shown by the net increase of the electron temperature in the plasma interaction region of few hundreds of eV. Time resolved X-ray images of the interaction region also provided valuable information on the dynamics of the interaction showing strong local heating of the plasma. This effect was found to be in good agreement with calculations performed assuming inverse bremsstrahlung absorption. These measurements also allowed an estimate of the electron thermal diffusion length, which was found to be consistent with the classical heat transport theory.

In summary, a preformed plasma with a good transverse electron density homogeneity was produced. The expansion fronts of the plasma were regular and it is predicted that, when irradiation takes place with an interaction beam focused in a spot of 100-200 µm with a large f/number, the interaction will start with negligible perturbations from refraction effects due to the limited amplitude of small scalelength density perturbations. These features, together with the temperature of several hundreds of eV at the time of interest, make these kind of plasmas extremely useful for studies devoted to a better understanding of the physics of the main instabilities affecting the corona of a laser irradiated micro-sphere in the inertial confinement fusion scheme. An experimental investigation on filamentation instability and Brillouin scattering, performed by interacting with preformed plasmas, like those described so far, will be presented in the next section.
4.6 - Filamentation Instability

An experimental study on laser interaction with underdense plasmas was carried out using a wide range of diagnostic techniques. Particular attention was devoted to the filamentation instability which, as discussed in Sect.2.3, can seed suitable conditions for other instabilities to grow. According to the conclusions of Sec.2.4, the onset of the filamentation instability can be related to second harmonic and X-ray emission processes. Experimental data concerning this aspect of laser-plasma interaction will be presented and discussed. The data shown here were obtained using the same experimental conditions described in Sect.4.1. In addition, some important results taken in a slightly different experimental configuration, nevertheless relevant to the topic under investigation, will also be presented. Time resolved X-ray images of the laser-plasma interaction region recorded employing the four-frame X-ray gated intensifier described in detail in Sec.4.3, will be shown which provide clear experimental evidence of the onset of whole beam self focusing. In addition, valuable information on the dynamics of the laser-plasma interaction was also gained from the interferometer described in Sect.4.2.

Filamentation instability and second harmonic emission

As already pointed out, the interferometric investigation described in Sect.4.2 allows second harmonic (SH) radiation emitted sidewards during the interaction process to be imaged out. As discussed in Sect.2.4, SH emission can take place during the heating of the plasma and during the interaction of the delayed pulse with the underdense plasma. It is therefore an important diagnostic tool to monitor the degree of uniformity of the plasma and, in particular, the onset of filamentation.

SH emission generated by the interaction beam in the underdense plasma was indeed detected by the interferometer in the experimental conditions described in Sect.4.2. Fig.4.6.1 shows an interferogram of the plasma taken 2 ns after the peak of the heating pulse at a total heating intensity was $9.4 \times 10^{13}$ W/cm$^2$. The interaction pulse was timed to reach the plasma 1.5 ns after the peak of the heating pulse and was focused on the plasma in a line-focus configuration with its longitudinal axis set along the line of view of the interferometer. The focal spot was
approximately 800 µm long and 100 µm wide and the average intensity was approximately $1.2 \times 10^{14} \text{ W/cm}^2$. Both fringe patterns produced by the interferometer, as explained in Sect.4.2, are displayed in Fig.4.6.1.

![Interferogram of a preformed Al plasma generated in similar conditions to those described in Sect.4.3, taken 2.0 ns after the peak of the heating pulses. The horizontal arrow on the RHS of the image indicates the direction of the laser beam while the position of the target in the double image produced by the interferogram (see Sect.4.2) is indicated by the vertical arrows.](image)

Bright localized SH emission sources arise from the interaction with a 600 ps pulse at an intensity of $1.2 \times 10^{14} \text{ W/cm}^2$ and delayed respect to the plasma forming pulse by 1.5 ns.

SH emission is generated by the interaction beam, as it propagates through the plasma, and is localised at the boundaries of the main bulk of the plasma. One can see that the SH radiation emitted at the output boundary is less intense, probably due to absorption of the interaction energy by the plasma. Incidentally we observe that the corresponding SH source in the two images produced by the interferometer have different intensities. According to the description of the interferometer given in Sect.4.2, the interferometer produces two orthogonally polarized images. The difference in the intensity of SH sources seen would therefore indicate that such SH radiation is partially polarized. Although further experimental investigation is needed to study the establish the importance of polarization effects, these circumstances could provide valuable insight into the dynamics of SH emission processes in underdense plasmas (Deha et al., 1992; Stamper et al., 1989).

On the other hand strong SH emission was also detected when the laser-heating action was localised onto a small portion of the target. **Fig.4.6.2** shows an interferogram obtained with the heating pulses turned off. The beam usually used for the delayed interaction was set to heat the Al target directly with the same line-focusing condition described above, at an average laser intensity of $1.0 \times 10^{14} \text{ W/cm}^2$. The probe pulse was delayed by 500 ps with respect to the main pulse. The strong SH emission at 90 degrees occurring in the proximity of the
target surface suggests that the intensity non-uniformities introduced by the focal spot boundaries and the consequent formation of density gradients in the plasma play a dominant role in the production of SH emission.

It is important to notice that, in contrast with the case of localized heating, no SH emission was detected when the dot Al target was uniformly irradiated in the centre of the laser focal spot as in the case of the interferograms of Figs.4.2.1-2. These circumstances strongly support the conclusion that the presence of transverse gradients in the laser intensity is a favourable condition for the generation of SH emission, both in the interaction with underdense plasmas and in the interaction with solid targets. Moreover, it is evident that, when conditions are suitable for SH generation, strong SH emission at 90 degrees to the incident laser axis is indeed detected and can be imaged using simple diagnostic techniques.

![Fig.4.6.2. Interferogram of plasma produced by a localised heating with a 600 ps, 1 µm laser beam of a 500 nm thick Al target, taken 0.5 ns after the peak of the laser pulse, showing strong SH emission in the proximity of the target surface. The horizontal arrow on the RHS of the image indicates the direction of the laser beam while the position of the target in the double image produced by the interferogram (see Sect.4.2) is indicated by the vertical arrows.](image)

Moreover, SH radiation emitted forward from laser-plasma filaments has also been detected using a gated optical intensifier. In this case a 3 ns, 1.064 µm laser beam was focused using an $f/8$ optics, in a 60 µm diameter focal spot on a 1 µm plastic (Formvar) foil at an average laser intensity in the focal spot of $3 \times 10^{13}$ W/cm$^2$. In this experiment (Gizzi et al., 1991) the maximum electron density at the peak of the 3 ns pulse was approximately a quarter of the critical density for the incident laser pulse and the electron temperature was (Gizzi et al., 1992) typically several hundred eV. **Fig.4.6.3** shows a time resolved image of SH emitted forward, i.e. in the direction of the incident laser beam. The object plane of the $f/8$ imaging optics was set at the target plane in order to image the laser-plasma interaction region. The gate time of the optical intensifier was 120 ps (FWHM) and the peak of the gate was located approximately at the peak of the
laser pulse. The filamentary regime taking place in these experimental conditions is shown by two important features. Firstly, the particular structure of the SH sources was found not to be reproducible shot by shot, as expected in the case that laser and/or electron density non-uniformities seed the conditions for the instability to occur. Secondly, the transverse size of the small features in the interaction region, typically ≈5 µm, is in very good agreement with the theory of filamentation as predicted by Eq.2.4.2, which, in the conditions of this experiments, gives a maximum spatial growth rate of ≈10 µm for a perturbation wavelength of 6 µm.

![Image](image_url)

**Fig. 4.6.3.** Time resolved image of the laser-plasma interaction region in the second harmonic (0.532 µm) of the interaction wavelength (1.064 µm), showing filamentary structures with ≈5 µm transverse size. The active temporal window of the gated optical intensifier used to record this image was 120 ps (FWHM) and was located at the peak of the 3 ns laser pulse.

In the case of Fig.4.6.3 the filamentation process involves portions of the incident beam and results in a breaking-up of the incident laser beam in several spots. However, in particular circumstances, plasma conditions may be such that the beam-plasma interaction region as a whole, becomes unstable leading to the so-called whole beam self-focusing as will be shown later in this chapter.

**Beam smoothing techniques**

Several possible solutions to the problem of filamentation, self-focusing and other harmful instabilities, have been suggested and their effectiveness is presently under investigation (Bradley et al., 1992; Afshar-rad et al., 1992; Peyser et al., 1991; Willi et al., 1990) and is one of the crucial issues in laser-plasma interaction
experiments. Particular attention is being devoted to techniques aimed to “smooth” the intensity distribution of the laser beam both spatially and spectrally, or, in other words, destroy the spatial and temporal coherence of the laser light. Random phasing (RP) (Kato et al., 1984), induced spatial incoherence (ISI) (Lehmberg & Obenschain, 1983) and smoothing by spectral dispersion (SSD) (Skupsky et al., 1989), are the most promising methods developed so far.

Long scalelength non-uniformities in the far field intensity distribution of a focused laser beam can be suppressed using the RP technique. In principle, the unfocused laser beam is divided into a large number of beamlets, typically hundreds to thousands, and each beamlet is then phase-shifted by an amount between 0 and $\pi$. In practice, however, due to the limitations imposed by manufacturing techniques, a binary-type phase plate array is typically used in which each beamlet is randomly phase-shifted by either 0 or $\pi$. The intensity distribution in the focal plane of a RP smoothed beam has an envelope function given by the square of the sinc function whose first zero contains 82% of the beam energy and is located at a distance from the beam axis of the order of $\lambda f / d$, where $\lambda$ is the laser wavelength, $f$ is the focal length of the focusing optics and $d$ is the size of each beamlet. However, due to interference among the beamlets, the far field intensity pattern is characterized by very strong stationary intensity fluctuations, whose transverse size is of the order of $\lambda f / D$, where $D$ is the beam aperture. These so called hot-spots can, in particular circumstances (Rose & DuBois, 1993), increase the likelihood of filamentation and/or beam self-focusing.

This problem can be partially solved in the case of broad-band lasers, by using ISI before performing the random phasing. In this case the beam is divided in beamlets, typically several tens, and a different delay is imposed on each beamlet by increasing its optical path. This can be operated using, for example, a set of two crossed “echelons”. The condition for this technique to work is that the delay increment must be larger than the longitudinal coherence time of the laser, $\tau_c = 1/\Delta \nu$. The far field intensity distribution resulting after focusing is a time-varying interference pattern which becomes smooth on a temporal scale of the order of several laser coherence times.

It has been pointed out (Skupsky et al., 1989) that the chaotic broad bandwidth required to implement ISI in a laser driver for inertial confinement fusion, would result in high-intensity temporal spikes in the laser pulse which could activate non-linear processes in the laser glass amplifier medium with consequent damage. Moreover, it would severely limit the efficiency of frequency up-conversion.
4.6 Filamentation Instability

Smoothing by spectral dispersion has therefore been suggested as an alternative to ISI. In the case of SSD a non-chaotic broad band-width laser light, achieved by means of electro-optical modulators, is used which can still be efficiently frequency tripled. The basic idea of SSD is to spectrally disperse such broad band-width light onto a random phase plate in order to irradiate each element of the plate with a different frequency. The interference pattern produced by the RPP in the focal plane will now rapidly vary in time on a time-scale of the order of $\tau_c$ and the temporally averaged intensity profile will converge to the square of the sinc function after this time, instead of several coherence times as in the case of ISI.

The effectiveness of RPP and ISI in reducing the level of instabilities in laser-plasma interactions has been confirmed by many experiments performed in various conditions (Willi et al., 1990; Obenschain et al., 1989, Coe et al., 1989). It is now clear that smoothing techniques can greatly reduce the development of density inhomogeneities which can arise from non-uniformities in the laser intensity on target. This smoothing effect is needed particularly in the early plasma formation process, when the plasma surrounding the solid target is too small for thermal conduction to be an efficient smoothing process.

However, uncertainty still remains (Kruer, 1991) on whether the onset of smoothing techniques is fast enough to provide a complete control over plasma inhomogeneities. With respect to this, recent experimental studies indicate (Desselberger et al., 1992) that non-uniformity imprinting on the target surface occurs very early in the laser pulse, typically in the first tens of picoseconds from the beginning of the ablation process. Furthermore these non-uniformities seem to persist during and after the laser pulse, even in the presence of laser beam smoothing. It has been suggested that magnetic fields can grow in the interaction region, early in the pulse due to density and temperature gradients, and establish themselves for the rest of the pulse, providing the conditions for density inhomogeneities to endure. This effect would seriously influence interaction processes as it would render smoothing techniques only marginally effective.

Instabilities driven by smoothed laser light

Experimental measurements have been performed in order to study the effect of induced spatial incoherence on laser interaction with long-scalelength plasmas. The preformed plasma was generated using the line-focus geometry described in detail in Sect.4.1. Stimulated Brillouin back-scattering (B-SBS) measurements were carried out simultaneously with time resolved X-ray imaging. The level of the B-SBS as a function of the intensity of the interaction beam is shown in Fig.4.6.4.
As described in Sect. 4.1, the preformed plasma was generated by irradiating Al stripe targets. Occasionally, as already discussed in detail in Sect. 4.4, KCl targets were also used. B-SBS data for both smoothed (crosses) and unsmoothed (open circles for Al plasmas and solid circles for KCl plasmas) laser pulses are plotted on the same graph. The configuration of the heating laser pulses was the same for all the data shown in Fig. 4.6.4 in order to produce the standard plasma conditions for the line-focus geometry already examined in Sect. 4.4. The electron density of the Al plasma, at the peak of the interaction pulse, was typically \(1.5 \times 10^{20} \text{ cm}^{-3}\) i.e. \(0.14 n_c\), \(n_c\) being the critical density for the 1 µm wavelength of the interaction laser, while the electron temperature was approximately 450 eV.

![Fig. 4.6.4. Stimulated Brillouin back-scattering level as a function of the intensity of the interaction beam with (crosses) and without (circles) induced spatial incoherence. The B-SBS level is given as a fraction of the incident laser energy. The data without ISI were recorded interacting with Al plasmas (open circles) and KCl plasmas (solid circles).](image)

According to these results, ISI seems to be very effective as it leads to a ten-fold reduction of the SBS-levels. In addition, similar results were found in the stimulated Raman back-scattering (Afshar-rad et al., 1992). However, when the conditions of the preformed plasma were modified in order to interact with a more collisional plasma, whole beam self-focusing occurred even in the presence of ISI and a ten-fold increase in the B-SBS level was detected. **Fig. 4.6.5** shows a sequence of X-ray images of the laser-plasma interaction region obtained in the same timing configuration of the sequence shown in Fig. 4.4.7 and Fig. 4.4.8. In this case the preformed plasma was produced by irradiating the 700 nm thick, 300 µm long, 300 µm wide Al target at an average heating intensity of
According to the numerical simulation, the maximum electron density of the preformed plasma, at the peak of the interaction pulse, i.e. 2.2 ns after the peak of the heating pulse, is expected to be approximately a factor of two higher than in the case of Fig.4.4.7-8, i.e. \( n_e \approx 3 \times 10^{20} \text{ cm}^{-3} \). Also, the electron temperature is predicted to be approximately 70% of the previous case, the new value being \( T_e \approx 300 \text{ eV} \). In contrast, the intensity of the interaction beam on the plasma was \( 2.5 \times 10^{14} \text{ W/cm}^2 \), i.e. almost a factor of four higher than in the case of Fig.4.4.7.

It should be noted however that, for the standard plasma conditions, no self-focusing was observed at the maximum intensity of the interaction beam available, that is, \( 8.7 \times 10^{14} \text{ W/cm}^2 \).

According to Eq.2.4.1, for a perturbation wavelength of the order of the interaction beam focal spot, i.e. 140 µm, and in the typical plasma conditions of interest in this study, the ratio between the effective thermal electron conductivity calculated solving the Fokker-Planck equation (as discussed in Sect.2.4), and the classical Spitzer-Härm conductivity is less than 0.3. In contrast, in the case of the standard preformed plasma conditions mentioned above, this ratio is approximately 0.15. Consequently, non-local effects in heat transport are expected to play an important role in determining the filamentation growth length for a given perturbation wavelength as discussed in Sect.2.4. The filamentation growth length given by Eq.2.4.2 has been evaluated for the standard plasma conditions of

\[ 1.6 \times 10^{13} \text{ W/cm}^2 \] on each side instead of the typical \( \approx 4 \times 10^{13} \text{ W/cm}^2 \).
Fig. 4.6.6 and for the conditions of Fig. 4.6.5 and is plotted as a function of the perturbation wavelength in Fig. 4.6.6.

According to these results, the increased collisionality of the plasma conditions relative to the event of Fig. 4.6.5 leads to an approximately 3-fold reduction of the filamentation growth length corresponding to the perturbation wavelength introduced by the focal spot of the interaction beam as a whole. In fact the growth length for an interaction intensity of $2.5 \times 10^{14}$ W/cm$^2$, for a perturbation wavelength of 140 µm, is 38 µm for the standard plasma conditions, i.e. $T_e \approx 450$ eV and $n_e \approx 0.14 n_c$ while it is only 13 µm for the conditions of the observed self-focusing, i.e. $T_e \approx 300$ eV and $n_e \approx 0.3 n_c$.

On the other hand, as already pointed out above, no self-focusing was observed in standard plasma conditions at the maximum interaction intensity available, that is, $8.7 \times 10^{14}$ W/cm$^2$. And in fact, the growth length corresponding to this interaction intensity is 20 µm, that is, less than a factor of two smaller than the 38 µm considered above, supporting the conclusion that the enhancement of collisionality was the primary cause of the observed occurrence of self-focusing. According to Fig. 4.6.5, and taking into account the angle of view of the X-ray imaging device (see Fig. 4.1.2), a self-focusing length of $\approx 280$ µm can be measured. A comparison between the values the growth-length given above requires a further discussion of the assumptions on which the current theory of filamentation itself, of which the model (Epperlein et al., 1990) used here is an extension (see Sect. 2.4.2), is based. The growth rate is (Berger et al., 1993) calculated in terms of the response
of the plasma to a small amplitude sinusoidal perturbation of the incident laser light. It is also assumed that the perturbation caused by the filamentation process to the background plasma, i.e. to the electron temperature and density, is small compared to their unperturbed values. As shown in Sect.4.5, this condition is not strictly fulfilled in our case as the interaction beam strongly perturbs the plasma during the interaction process even when no filamentation occurs. In addition, the attenuation of the laser light, due to absorption by inverse bremsstrahlung, is not taken into account.

Therefore, although these models give the likelihood of having favourable conditions for thermal or ponderomotive self-focusing to occur, they tend to overestimate the spatial growth rate. An estimate of the relevance of these effects can be obtained from a simplified model (Estabrook et al., 1985) in which the self-focusing length is determined according to the classical filamentation theory with attenuation effects treated self-consistently. For our standard experimental conditions this model predicts a 25-fold increase of the self-focusing length compared to the classical value. Furthermore, this factor is weakly dependent on the plasma conditions. If we assume a similar trend for the value calculated using the non-local theory we obtain a self-focusing length of 1mm, which is approximately equal to the plasma longitudinal extent. In contrast, using the same procedure to evaluate the growth rate relative to the conditions for which self-focusing was observed, we find a self-focusing length of approximately 300 µm which is consistent with the measured value of 280 µm.

In summary, the measurements presented in this Section provide a good test-bed for the present models of the filamentation instability. The results discussed here were found to be consistent with the classical theory of filamentation provided that non-local electron transport effects are taken into account and that the attenuation of the interaction beam by inverse bremsstrahlung absorption is also included in the calculation of the spatial growth length of the filamentation instability.
CHAPTER 5

SHORT-SCALELENGTH PLASMA STUDIES

An experimental investigation on laser-plasma interaction, in a picosecond regime, will be presented in this Chapter. Short-scalelength plasmas were produced by irradiating solid targets using high intensity picosecond UV laser pulses. Absorption measurements showing the effect of polarization and angle of incidence will be presented, followed by a detailed numerical modelling. Measurements of transient heat propagation in solid matter, carried out using picosecond time-resolved X-UV spectroscopy, will be described. The interplay between absorption processes and heat transport phenomena will be analysed in detail, by using simple analytical models as well as the numerical simulations already presented in this thesis.
5.1 - Picosecond Laser–Plasma Interaction

An experimental investigation on short-pulse laser-produced plasmas will be described in this Chapter. Results will be presented on the interaction of picosecond laser pulses with matter in a high intensity regime (>10^{16} \text{ W/cm}^2). Short scalelength plasmas were produced irradiating solid targets with a prepulse-free, 12 ps, 0.25 µm laser pulses. Absorption measurements will be discussed which provide information on the interplay between laser-plasma energy transfer mechanisms and electron heat transport. These results are related to X-UV spectroscopic measurements with picosecond temporal resolution which enabled the first observation of transient heat propagation in solid targets.

A detailed description of the experimental set-up employed in the investigation will be described in this section with particular attention to the features of the laser source which are relevant to the experimental investigation presented in the following sections. The absorption experiment will be presented in Sect.5.2 and discussed in Sect.5.3. X-ray self emission processes are considered in Sect.5.4 with emphasis to those results related to the characterization of plasmas in terms of electron density and temperature. Finally, heat transport phenomena are studied and discussed in detail in Sect.5.5.

Features of the laser source

The experimental investigation was carried out at the Central Laser Facility of the Rutherford Appleton Laboratory, using the SPRITE laser system (Ross et al., 1990). The KrF pumped Raman-shifted laser operated at a wavelength of 268 nm and delivered up to 6 J in a 12 ps FWHM pulse. The intensity ratio between the short pulse and the long pulse background, due to amplified spontaneous emission (ASE), was measured (Harvey et al., 1991) to be greater than 2\times10^{10}, the ASE emission being 10 ns in duration. The 8 cm diameter flat-top laser beam was focused on target using a reflecting optics consisting of an F/3 off-axis parabolic mirror (Ross, 1991). Reflective optics were adopted in this high intensity regime, in place of the transmitting optics in order to avoid energy losses due to two photon absorption (Liu et al., 1978). In fact in our experimental conditions the maximum
incident irradiance on the focusing optics was typically $\approx 5$ GW/cm$^2$ and therefore non-linear mechanisms could take place in transmitting media. The effect due to two photon absorption for e.m. waves propagating through a non-linear medium is described by the following equation

\[ d I(x) = -\left[aI(x) + bI(x)^2\right] dx, \quad (5.1.1) \]

where $I(x)$ is the intensity of light after it has propagated over a distance $x$, $a$ is the linear absorption coefficient and $b$ is the non-linear absorption coefficient.

The absorption coefficient accounting for the two photon effects in quartz at a wavelength of 266 nm has been measured (Liu et al., 1978) to be $b = 4.5 \times 10^{-2}$ cm / GW. Eq.5.1.1 was solved numerically assuming non-linear absorption only. The fractional loss of laser intensity as a function of the incident intensity is shown in Fig.5.1.1 for propagation through a 1 cm (solid line) and 2 cm (dashed line) thick quartz plate. According to this result, at the maximum available irradiance two-photon absorption becomes efficient and can lead to an absorption of as much as 40% of the laser energy.

![Fig.5.1.1.](image)

Fig.5.1.1. Transmittivity of a 1 cm (solid line) and 2 cm thick (dashed line) quartz plate, as a function of the intensity of UV (0.25 µm) light assuming two-photon absorption only.

A multilayer reflecting optics can easily exhibit a 99% reflectivity free from non-linear absorption effects due to the small thickness of the layers of the coating. The off-axis parabolic mirror adopted in our experiment had a reflectivity greater than 99% at both the Raman shifted wavelength (266nm) and the unshifted KrF wavelength (248nm). The use of other UV transmitting materials characterized by a lower two photon absorption coefficients can also be considered in order to reduce
non-linear effects. In the case of LiF or CaF$_2$, for example, the non-linear absorption coefficient was found (Liu et al, 1978) to be $< 2 \times 10^{-2}$ cm / GW in both cases. On the other hand, these materials are highly hygroscopic and therefore tend to degrade very soon. For this reason they cannot be used as lens media and were only employed on the laser beam input window of the interaction chamber.

**Experimental set-up**

The configuration of the laser beam is shown in Fig.5.1.2 for the case of the absorption measurements described in Sects.5.1-2. Besides the particular arrangement of the integrating sphere necessary for the absorption measurements, the remaining part of the experimental set-up was adopted throughout the experimental investigation described in this Chapter. Also shown are the main diagnostics employed to monitor the laser performance on a shot-by-shot base. According to the experimental set up shown in Fig.5.1.2, taking into account the losses due to optical components on the beam line, the total transmittivity from the laser output to the target was approximately 75%. This value, required to estimate the energy incident on target, was also confirmed by accurate beam calorimetry.

The beam was then focused on target with an effective $f / 4.7$ aperture off-axis parabolic mirror. According to the diffraction theory for uniform illumination of a focusing optics with circular aperture and $f$ - number $F$, the depth of focus $z$ and the focal spot diameter $d$ are respectively

\[
    z = \pm 2 \lambda \cdot F^2 \theta / \theta_{DL} \tag{5.1.2}
\]

\[
    d = 2.44 \lambda \cdot F \theta / \theta_{DL} \tag{5.1.3}
\]

where $\theta$ is the actual divergence of the beam and $\theta_{DL}$ is the diffraction limited beam divergence. The beam divergence in our experimental conditions was measured (Harvey et al., 1991) to be approximately five times the diffraction limited value, resulting in a focal spot diameter, at the best focus, of $\approx 15$ µm and a depth of focus ±50 µm.

According to the particular experimental condition required, the intensity on target was also varied by changing the size of the focal spot on target. This was accomplished by moving the target towards the focusing optics. The size of the focal spot was measured from comparison with a solid sample, typically a thin wire with a known diameter, using an *obscuration* technique. The value given hereafter of the focal spot diameter on target corresponds to the diameter of a wire that, placed
at the target plane, would completely obstruct the beam. Therefore, the diameter of the wire is equal to (and not greater than) an effective transverse size of the beam. This procedure was carried out with the laser running in a regime of very low energy per pulse and with a repetition rate of several pulses per second. The obscuration was determined by monitoring the shadow of the wire in the beam, on fluorescent paper, at a distance after the focal plane comparable with the focal length of the focusing optics, and by varying the distance between the optics and the target.

![Fig.5.1.2](image)

**Fig.5.1.2.** Experimental set-up showing laser beam configuration and diagnostic arrangement for the shot-by-shot monitoring of the laser performance. Also shown is the arrangement for the absorption measurements described in detail in the following sections.

The maximum intensity available on target, at the best focus, was therefore expected as high as $10^{17}$ W/cm$^2$ while the intensity due to the ASE was less than $10^7$ W/cm$^2$. According to recent studies on laser induced damage threshold (Corkum et al., 1988), substantial plasma formation in laser-solid interaction can occur if the irradiance exceeds $10^9$ W/cm$^2$. Consequently, the ASE prepulse present in the experimental conditions described here is not expected to generate sizeable plasma and the 12 ps pulse is expected to interact directly with the solid.

Experimental evidence of the absence of preformed plasma was obtained from measurements performed by irradiating Aluminium targets coated with a very thin plastic coating. A full description of these measurements will be given in Sect.5.5. X-UV time resolved spectroscopy with picosecond resolution was employed to study the propagation of the heat front in such targets at irradiances between $10^{16}$ and $10^{17}$ W/cm$^2$. The temporal evolution of X-UV line emission clearly indicated that the interaction of the ASE pulse with a layer of plastic as thin as 1000 Å resulted in negligible plasma effects.
The intensity distribution in the focal plane was experimentally monitored using an equivalent plane (EP) system. The fraction of the laser beam transmitted by an high reflectivity wedged mirror placed on the main beam-line was focused using an 8 m focal-length quartz lens to obtain a 24x magnified image of the focal spot. The far field pattern recorded by the EP monitor was stable shot by shot and, at the best focus, showed typically a 20 µm focal spot diameter which is consistent with the 15 µm calculated above according to Eq.5.1.3.

An X-ray pin-hole camera (see Sect.4.4) fitted with a 5 µm diameter pin-hole and filtered with a 20 µm thick Beryllium foil was employed to obtain images of the plasma in the keV spectral region with a magnification of 16x. In this case the limit to the image resolution, set by the fibre optics of the phosphor plate, was approximately 100 µm which corresponded to ≈7-8µm in the plasma. Occasionally a higher (30x) magnification was used in order to resolve details of the order of the pin-hole diameter (≈5µm) in the plasma. Under conditions of best focusing, the pin-hole camera showed a basically circular image of 20 µm FWHM which is consistent with the calculation and the equivalent plane images.
5.2 - Techniques for Absorption Measurements

When a high power laser beam interacts with a solid target, laser energy is partially absorbed by the plasma and partially scattered in a broad (≈2π sterad) angle with the angular distribution being determined by the particular interaction regime taking place as well as by the angle of incidence of the laser beam. Measurements of the energy absorbed by the plasma can therefore be performed by collecting the fraction of laser light scattered by the target during the interaction.

Experimental set up for absorption measurements

From an experimental viewpoint we distinguish between collimated scattering and diffuse scattering. Collimated scattering of incident laser light can originate from reflection of the laser light at the critical density surface. However, before reaching the critical surface, the laser light propagates in an underdense plasma where laser induced instabilities can also be activated. According to the conclusions of Sect.2.3, laser driven instabilities can efficiently grow provided that the density scalelength of the plasma in which the laser light is propagated is sufficiently large. However, despite the short density scalelength generated by short pulse laser produced plasmas, the stimulated Brillouin scattering described in Sect.2.3 can be, particularly at high laser irradiance, an effective scattering mechanism.

The ideal detector to employ in measurements of scattered laser light in laser-plasma interaction experiments should therefore have a wide enough collection angle in order to detect both components of scattered light. A 10 cm diameter Ulbricht sphere (Godwin et al., 1977) was employed in our experiment in order to collect the light scattered by the target placed at the centre of the sphere. The inner surface of the sphere was coated with a Barium Sulphate diffuse reflecting paint whose reflectivity at 268 nm is better than 90%. The high reflectivity allows each ray of the laser light scattered by the target/plasma to undergo many reflections and eventually escapes from a small aperture of the sphere where a detector is placed.

This process is illustrated in Fig.5.2.1 where the set up of the sphere and the detector is schematically shown. Since the sphere has a high collection efficiency over a wide spectral range from 250 nm up to 1100 nm, plasma self-emission in this
spectral range is also collected by the sphere and contributes to the measured scattering. However, typically a small fraction of the order of 10% of the laser energy transferred to the plasma is converted into self-emission and mostly in the X-ray spectral range (<10 nm) which cannot be detected. Consequently, contributions from self-emission to the collected radiation energy can be neglected.

A 1cm diameter hole on the sphere allowed the laser beam to be focused onto a solid Al target with an \( f/4.7 \) effective \( f \)-number. The light scattered by the target and collected by the sphere was detected using a Gentec calorimeter placed on the sphere output hole. A back-scattering channel was also set in order to detect the fraction of light scattered by the target back into the beam input hole. Three 5 mm diameter diagnostic holes of the sphere were used for the X-ray pin-hole camera, for X-ray imaging, and target injection and positioning. The size and the number of these apertures on the sphere was kept as small as possible to minimize losses and consequent reduction of collection efficiency, provided that a satisfactory monitoring of the interaction conditions could be carried out shot by shot.

![Experimental set up of the Ulbricht sphere and the back-scattering channel for collection of the light scattered by the solid Aluminium target for S and P polarized laser light.](image)

On the other hand, these apertures do not affect the linear relation between scattered energy and output energy unless they intercept part of the collimated scattering before any diffuse reflection has occurred. In fact it can be shown that the fraction of light diffused by a point on the sphere and going through a small diagnostic hole is independent of the position of the hole itself with respect to the diffusion centre and is given by the ratio between the area of the hole and the area of the sphere surface. Due to geometrical considerations, only 0.015% of the light diffused by any point of the inner surface of the sphere is lost through a 5 mm diameter hole, resulting in a depletion of the sphere collection efficiency. Provided that the reflecting coating behaves linearly with the incident energy flux, the
collection efficiency is a function of the reflectivity of the coating and of geometrical factors only. A calculation of the overall collection efficiency of the sphere for a given configuration is though a very difficult task and the poor knowledge of the response of the coating to high power laser light required a careful calibration process to be carried out.

Both the sphere and the detector employed to detect the collected light were calibrated before and after the measurements. The calibration of the diffuse scattering channel was performed by replacing the target at the centre of the sphere with a beam block coated with diffuse reflector and set in a grazing incident configuration. The incident laser beam was defocused in order to bring the irradiance on the block well below the damage threshold. A focal spot diameter of ≈1cm was used, giving a maximum irradiance below ≈10^{10} W/cm^2. In this way ≈100% of the incident laser light was scattered in a diffuse manner into the sphere allowing the overall collecting efficiency of the sphere to be measured. Measurements taken in a laser energy regime from 100 mJ to 4 J consistently gave a sphere collecting efficiency of 0.02 ± 0.001.

A plane 25 µm thick Al target mounted on a remotely driven translatable rotating support was placed in the centre of the integrating sphere. Since the laser light was already linearly polarized in a vertical direction, S and P polarizations were simply obtained by changing the axis of target rotation. A rotation of the target around a vertical axis resulted in an S-polarized incidence while a rotation around an horizontal axis normal to the beam resulted in a P-polarized incidence. Fig.5.2.2 illustrates these two configurations.

![Fig.5.2.2](image)

**Fig.5.2.2.** Schematic arrangement of laser-target configuration for S and P polarized incidence. The axis of target rotation was horizontal for S-polarization (side-view on the left) and vertical for P-polarization (top-view on the right).

A high magnification optical imaging system looking at the target through a small diagnostic hole of the sphere allowed the target position to be controlled to within less than 10 µm along the laser beam axis. Since the depth of focus in our
configuration (see Sect.5.1) was ± 50 µm, the shot to shot fluctuations of incident laser irradiance due to fluctuations of the focal spot size were very small. The focal spot size was also monitored using an X-ray pin-hole camera. The size of the X-ray emitting region was found to be reproducible shot by shot.

Absorption as a function of laser intensity

The dependence of the absorption as a function of the laser intensity was carried out in a normal incidence configuration. The intensity on target was varied by defocusing the laser beam from the 20 µm diameter focal spot up to a 150 µm focal spot. Fig.5.2.3 shows the dependence of total scattered laser energy as a function of laser intensity on target in the range between a few times $10^{14}$ and $5 \times 10^{17}$ W/cm$^2$. The fraction of scattered laser light increases with intensity from approximately 40% at $10^{14}$ W/cm$^2$ to more than 70% at the maximum intensity.

![Graph showing absorption as a function of laser intensity](image)

**Fig.5.2.3.** Experimental absorption as a function of incident irradiance on target at normal incidence. The vertical error bars are given by the uncertainty on the collection efficiency of the sphere while the horizontal ones derive from the accuracy of target positioning as well as on the uncertainty of laser energy measurements.

According to this graph, the absorption of laser energy by the plasma decreases rapidly with the laser intensity. A detailed modelling (Riley et al., 1993) of absorption as a function of the incident intensity based on the hydrodynamic numerical code MEDUSA, shows that the observed behaviour is basically in agreement with the calculations. In particular, for irradiances below $10^{16}$ W/cm$^2$, inverse bremsstrahlung is found to account for most of the absorbed laser energy.
Beyond this point calculations indicate that resonance absorption becomes the dominant absorption mechanism. In view of the fact that even at this high irradiance on target no early (ASE induced) plasma formation takes place, it is interesting to study absorption processes as a function of laser polarization and angle of incidence. The experimental results relevant to this aspect, presented and fully analyzed in the next section, will allow us to clarify the role played, in the laser-plasma interaction regime under investigation, by the two most important absorption mechanisms, namely inverse bremsstrahlung and resonance absorption.
5.3 - Polarization Effects on Absorption

The effect of the polarization of the incident laser light on laser-plasma coupling mechanisms was investigated by measuring the angular dependence of the absorption coefficient for both S and P-polarized light, using the experimental set up described in Sect.5.2. The experimental results obtained from this study will be presented in this section. A detailed modelling of these results will be given with special attention on the role played by non-local heat transport phenomena and non-linear effects in the inverse bremsstrahlung process.

Experimental measurements

Several experiments (Maaswinkel et al., 1979; Balmer, et al., 1977) have investigated the coupling mechanisms of short laser pulses with solid matter. The role of laser light polarization on the absorption of several tens of picoseconds pulses, typically between 30 and 50ps, incident on solid plastic targets was studied for 1μm and 0.53μm laser light. These measurements established the effectiveness of resonance absorption in the interaction of short pulse P-polarized laser light with small scalelength hot plasmas as predicted by theory (Pearlmann & Matzen, 1977).

More recent experiments (Kieffer et al., 1991; Fedosejevs et al., 1990) have extended this study to the regime of ultra-short (≤1 ps) laser pulses and shorter wavelength, down to 248 nm. In this regime the contribution of hydrodynamic phenomena is negligible and the interaction can be described in terms of linear metal optics theory. However, the effect of ASE was found to strongly affect the interaction processes in the high intensity regime, due to early plasma formation.

In the experiment described here solid Al targets were irradiated with prepulse-free, 12 ps, 0.268 μm laser light at an average intensity $5 \times 10^{16}$ W/cm$^2$ in a ≈ 20 μm diameter focal spot. According to the predictions of the hydrodynamic numerical simulation discussed in Sec.3.1, the interaction is expected to generate a plasma that, at the peak of the laser pulse, is characterized by a longitudinal density scalelength at the critical density of a few microns with an electron temperature in the sub-critical region of the order of few keV. In order to understand the interplay between the various processes which can account, in this particular regime, for transfer of laser energy to the plasma, we studied the dependence of the absorption
coefficient upon the angle of incidence and the polarization of the laser light. Since
the absorption measurements were time-integrated, the information obtained will
somehow be averaged over the whole interaction process. In the particular case of
absorption measurements, however, simple energetic considerations suggest that the
interaction processes occurring in a fraction of the pulse-length around the peak of
the pulse are expected to give a dominant contribution. This aspect will be
accurately discussed below.

**Fig.5.3.1** shows the dependence of the fractional scattering of laser light upon
the angle of incidence in the case of S-polarized light. The configuration of the laser
E-field relative to the target is schematically shown in Fig.5.2.2. A maximum
absorption of approximately 30% occurs at normal incidence. The absorption
decreases with increasing angle of incidence and almost 100% of the laser light is
scattered by the target for angles of incidence greater than 40 degrees.

![Fractional scattering of laser energy incident on a solid Al target](image)

**Fig.5.3.1.** Fractional scattering of laser energy incident on a solid Al target at an intensity of
5×10^{16} W/cm\(^2\) as a function of the angle of incidence and for S-polarized light. Both diffused
scattering and collimated back-scattering were collected and measured. The solid and dashed lines
show the result of the calculation (see text).

According to the theory of propagation of S-polarized e.m.waves in
inhomogeneous plasmas, the primary energy transfer mechanism is collisional or
inverse bremsstrahlung absorption. The basic features of these processes have been
discussed in Sec.1.2, 1.3 and 3.2. According to that analysis the absorption
efficiency at a given angle of incidence depends primarily upon two plasma
parameters, namely the electron density scalelength and the electron temperature of
the sub-critical density region. A relatively precise measurement of the density
scalelength is provided by the absorption measurements as a function of the angle of
incidence for P-polarized light. In fact, as discussed in Sect.3.2, a maximum
absorption is expected for a given angle of incidence at which maximum conversion of laser energy into longitudinal electron plasma waves occurs. According to Eq.3.2.3, this angle does not depend upon the electron temperature, being a function of the density scalelength and the wave-number of the laser light only. Therefore, a relatively direct measurement of the density scalelength can be obtained from the measurement of the angle of maximum absorption.

The dependence of the scattered light measured using the same condition of Fig.5.3.2, but with P-polarized light, is shown in Fig.5.3.2. The minimum fractional scattering and therefore, the maximum absorption, occurs at an angle of incidence $\theta_{\text{max}} \equiv 10 \pm 1\,\text{deg}$. At this angle approximately 45% of the incident laser energy is transferred to the plasma, partially via collisional absorption and partially via resonance absorption as shown in the plot of Fig.3.2.3. According to Eq.3.2.3 the measured angle of maximum absorption gives a scalelength of $L = 4.3 \pm 1.2\,\mu\text{m}$.

![Fig.5.3.2.](image)

Fig.5.3.2. Fractional scattering of P-polarized laser light incident onto a solid Al target at an intensity of $5 \times 10^{16}\,\text{W/cm}^2$ as a function of angle of incidence. Both diffused scattering and collimated back-scattering were collected and measured.

It is instructive to compare this result with a simple estimate of the density scalelength obtained by considering a self-similar expansion of a one-dimensional isothermal plasma (Kruer, 1988). According to this model the density scalelength increases linearly with the time $t$, and is simply given by

$$L = v_s t,$$  \hspace{1cm} (5.3.1)

where $v_s$ is the speed of sound given by Eq.1.2.3 which, in the case of a fully ionized Al plasma at an electron temperature of 2 keV is $0.35\,\mu\text{m/ps}$. Therefore, after approximately 10 ps from the beginning of the interaction the density
5.3 Polarization Effects on Absorption

scalelength is predicted to be 3.5 µm which is in agreement with the measured value reported above. A more accurate modelling of the temporal evolution of the density scalelength at the critical density was obtained from the hydro-code MEDUSA described in Sect.2.1 and Sect.3.1 and is reported below.

**Numerical hydrodynamic modelling**

The electron density profile of the plasma, at the peak of the laser pulse, for the experimental conditions given above in this section, is given in Fig.3.1.2. The portion of the profile for densities just below the critical density was fitted using a linear function and the corresponding scalelength was calculated according to its definition, \[ L = n_e / (\partial n_e / \partial x) \]. The same calculation was performed for analogous plots taken at different times throughout the interaction. The result is shown in Fig.5.3.3 as a function of the time relative to the peak of the pulse.

![Graph showing electron density scalelength at the critical density as a function of time](image)

**Fig.5.3.3.** Electron density scalelength at the critical density as a function of the time with respect to the peak of the 12 ps laser pulse as calculated by MEDUSA. The 268 nm laser pulse-shape was assumed to be Gaussian and the intensity on the solid Al target was \( 5 \times 10^{16} \text{ W/cm}^2 \).

According to this result, two regimes of expansion can be identified, one before and one after the peak of the laser pulse respectively. In the first stage of expansion before the peak of the pulse, the scalelength increases in time with a rate of approximately 0.3 µm/ps, which is consistent with the simple estimate of the speed of sound given above. In particular, the predicted scalelength at the peak of the pulse is 4.1 µm which is also in good agreement with the measured value of \( L = 4.3 \pm 1.2 \) µm. However, after the peak of the pulse, a rapid increase in the expansion velocity is predicted by the 1-D simulation, the rate being now 1 µm/ps,
that is, almost three times larger. The increase of the electron temperature occurring at the peak of the pulse and the consequent increase of the speed of sound is the main reason for this sudden change in the hydrodynamic expansion velocity.

A more accurate comparison between the experiment and the prediction of the hydrodynamic simulation requires that a time-averaged value of the calculated scalelength is taken into account, the average being weighted according to the intensity of the laser pulse. A fifth order polynomial fit of the curve of Fig.5.3.3 was computed and the weighted average scalelength, \( \bar{L} \) was calculated according to the usual definition

\[
\bar{L} = \frac{\Delta t}{\int_{-\Delta t}^{\Delta t} L(t) I_N(t) dt},
\]

where \( I_N \) is the gaussian laser intensity normalized according to \( \int_{-\infty}^{\infty} I_N(t) dt = 1 \), and \( \Delta t \) is large compared to the FWHM of the laser pulse. This calculation gives an average scalelength \( \bar{L} = 5.5 \, \mu m \). Although this value is still consistent with the experimental one once the error is taken into account, nevertheless it suggests that the simulation tends to overestimate the scalelength. This effect is expected to be particularly important after the peak of the pulse, when the longitudinal density scalelength becomes comparable with the transverse scalelength, that is, of the order of the focal spot radius, i.e. \( \approx 10 \, \mu m \). In these circumstances 2-D effects start to contribute to the hydrodynamic expansion of the plasma. As a consequence of the 1D expansion compared to the real expansion, simple thermodynamic considerations lead to the conclusion that the 1-D calculation is likely to overestimate the longitudinal density scalelength as well as electron temperature in the sub-critical region. On the other hand, as shown in Fig.3.2.1, the electron temperature in the underdense region strongly affects the collisional absorption since, according to Eq.1.2.13, the inverse bremsstrahlung absorption coefficient depends upon \( T_e^{-3/2} \) once the weak dependence of the Coulomb logarithm upon the electron temperature is neglected.

**Modelling of S-polarized absorption data**

A modelling of the absorption of S polarized light as a function of the angle of incidence can now be performed assuming a linear density profile with a scalelength given by \( L = 4.3 \pm 1.2 \, \mu m \). Some resonance absorption can take place in this case due to the spreading of angle of incidence in the \( f/4.7 \) focusing optics. However, we notice here that with the measured value of the scalelength, our plasma is entirely
5.3 Polarization Effects on Absorption

confined in the depth of focus of ±50µm given in Sect.5.2. In this region the laser beam wave-front is approximately flat and the incident light can be considered a plane wave over the whole longitudinal plasma extent. Some residual spreading can still arise from laser beam non-uniformities which can result in a local distortion of the wave-front. However, such contributions are expected to be effective only at small angles of incidence, typically of the order of the angular semi-aperture of the focusing optics. Therefore, absorption of S-polarized light for $\theta \leq 6$ deg should be mostly accounted for by collisional absorption.

The expected contribution of collisional absorption has been evaluated using the model discussed in Sect.3.2, assuming propagation of the laser beam in a stratified plasma with an exponential density profile. Refraction effects on the laser light reflected at the turning point and propagating backwards down the density gradient have not been taken into account in this analysis. In fact, there is experimental evidence that most of the laser light not absorbed by the target was scattered in an approximately $f/1.5$ cone. Fig.3.2.1 shows the fractional scattering calculated according to Eq.3.2.2 for a scalelength of 4 µm and for the electron temperature ranging from 1 to 5 keV. From a comparison of these plots with the experimental data of Fig.5.3.1, good agreement is obtained if the curve at the electron temperature of 5 keV is considered. Incidentally, this value of the electron temperature is also in good agreement with the prediction of the numerical simulation, although in this case the comparison with the simulation is limited to the normal incidence case. In fact, according to the plot of Fig.3.1.1, the electron temperature is very much uniform over the whole subcritical region, its value being approximately 6 keV. However, in view of the conclusions formulated above on the validity of the 1D simulation, the apparent consistency between experiment and calculation could, in this case, be only fortuitous and the simulation results must be carefully thought about in these circumstances.

Subcritical plasma and non-local heat transport

Although there are no direct experimental measurements of the electron temperature in the underdense region, valuable information can be gained from measurements concerning the supercritical region. Detailed measurements of the temporal evolution of the electron temperature have indeed been performed (Riley et al., 1992) using time resolved X-ray spectroscopy of K-shell Al line emission from H-like and He-like ions. This study was mainly concerned with the high density plasmas, as most of the observed emission originated from the supercritical region.
According to these measurements, at an intensity on target of $6 \times 10^{15}$ W/cm$^2$, an electron temperature up to 1 keV could be inferred, at the peak of the pulse, from comparison of experimental line ratios with the predictions of a time dependent atomic physics simulation. On the other hand, Stark broadening of these lines also indicates that the observed X-ray emission was generated at densities of several times the critical density. Therefore, the measured electron temperature should be related to this plasma density region.

On the other hand, recent (Rickard et al., 1989) 2-D Fokker-Planck simulations performed to study the interaction of 3.5 ps, 248 nm laser pulses with solid Al targets provide valuable information on the plasma condition in the subcritical region. The interaction of a 20 µm diameter focal spot with an Al plasma at an analogous intensity on target as that considered above, that is, $6 \times 10^{15}$ W/cm$^2$, was considered in that work. While predicting an electron temperature of 1 keV at two times the critical density, this study shows that the electron temperature in the underdense region is approximately 1.5 keV. Moreover, it also shows that, although lateral heat flow is negligible in the supercritical region, it plays an important role in the underdense region where free streaming lateral heat flow occurs. In other words uninhibited lateral heat flow is expected to occur in the underdense region once focal spot diameter becomes as small as few tens of the laser wavelength as in the case of this study.

There is consequently strong evidence that, at a laser intensity greater than $10^{15}$ W/cm$^2$, there is an interaction regime in which further increase of incident laser energy flux has little effect on the electron temperature in the underdense plasma, as thermal energy can escape efficiently from the interaction region. Since these effects, as well as 2D expansion effects, are not taken into account for by the 1-D hydrocode, it is expected that the electron temperature predicted by MEDUSA is largely overestimated.

**Laser field swelling and non-linear effects**

Another important aspect of the regime under investigation is that, at the irradiances typical of this experiment, the electron quiver velocity, $v_q$, is comparable to the electron thermal velocity, $v_{th}$. At the laser intensity of $5 \times 10^{16}$ W/cm$^2$, the field strength parameter introduced in Sec.3.3, i.e. the ratio between the quiver and the thermal velocity, is 0.44 in the case of a 2 keV plasma and becomes 0.62 for a 1 keV plasma. In addition, according to Eq.1.3.3, the swelling of the laser field in the plasma density gradient can give up to a three-fold increase of this parameter.
Therefore, in the proximity of the critical layer the quiver velocity can become higher than the local thermal velocity and consequently lead to a transition of the inverse bremsstrahlung absorption to a non-linear regime.

The contribution of non-linear effects in the collisional absorption of S-polarized light has been evaluated solving the differential equation given by Eq.3.3.5 with a linear density profile, assuming an intensity dependent absorption coefficient and including swelling of the E-field in the density gradient. Consistent with the conclusion reached above on the electron temperature in the subcritical region, an indicative temperature of 2 keV was assumed in the calculation.

The contribution of the two non-linear mechanisms discussed in Sec.3.3, i.e. the Langdon effect and the standard non-linear effect have been analysed separately. In the case of the Langdon effect, the non-linear contribution was introduced in the analysis by using the corrected absorption coefficient given by Eq.3.3.2 while the standard non-linear effect was taken into account using the absorption coefficient given by Eq.3.3.4. The swelling of the E-field in the density gradient was accounted for by replacing the vacuum E-field of the light with that given by Eq.1.3.1. The results of this calculation are summarised in Table 5.3.1 for the case of normal incidence, and for different values of the incident intensity ranging from $10^{13}$ W/cm$^2$ to $10^{17}$ W/cm$^2$. The intensity at the turning point, and at the output of the 4 µm density scalelength plasma, after reflection, is given as a fraction of the incident intensity, expressed in units of $10^{17}$ W/cm$^2$.

<table>
<thead>
<tr>
<th>$I_{in}$</th>
<th>$I_{cr}^{Lang}$</th>
<th>$I_{cr}^{snl}$</th>
<th>$I_{out}^{Lang}$</th>
<th>$I_{out}^{snl}$</th>
<th>$A_{Lang}(%)$</th>
<th>$A_{snl}(%)$</th>
</tr>
</thead>
<tbody>
<tr>
<td>1.0E-4</td>
<td>5.1E-5</td>
<td>5.1E-5</td>
<td>2.9E-5</td>
<td>2.7E-4</td>
<td>71</td>
<td>73</td>
</tr>
<tr>
<td>0.1</td>
<td>0.07</td>
<td>0.067</td>
<td>0.05</td>
<td>0.044</td>
<td>50</td>
<td>56</td>
</tr>
<tr>
<td>0.5</td>
<td>0.36</td>
<td>0.43</td>
<td>0.27</td>
<td>0.37</td>
<td>46</td>
<td>27</td>
</tr>
<tr>
<td>1.0</td>
<td>0.73</td>
<td>0.93</td>
<td>0.55</td>
<td>0.86</td>
<td>45</td>
<td>14</td>
</tr>
</tbody>
</table>

Table 5.3.1. Absorption of a 268 nm laser light after propagation and reflection in a 2 keV planar inhomogeneous plasma with a linear density profile of 4 µm scalelength. The incident intensity in the first column on the left is in units of $10^{17}$ W/cm$^2$ while the absorption coefficient in the last two columns on the right is given in percentage of the incident intensity.

In the limit of low intensity, reported in the first row of the table, non-linear effects are expected to be negligible and in fact the calculation is in agreement with the value given by Eq.3.2.2 and Fig.3.2.1 for normal incidence and a 2 keV electron temperature. At an intensity of $5 \times 10^{16}$ W/cm$^2$ (third row from the top), i.e. in the conditions of the absorption experiment described here, the calculation shows a strong reduction of the absorption due to non-linear effects. In particular, the
standard non-linear effect by itself reduces the absorption efficiency from the 73% of the linear case to 27% which is in agreement with the measured value given in Fig.5.3.1. In this particular case the calculation was performed for different angles of incidence in order to directly compare the calculation with the experimental results and the result is summarised by the solid line of Fig.5.3.1. In addition the dashed curve in the same graph shows the result obtained using the same condition but without the swelling effect.

The striking agreement between the experimental data and the results of this simple model indicates that, in the interaction conditions investigated, non-linear effects in the collisional absorption of laser light can play a very important role. On the other hand, it is important to point out that the extent to which these effects influence the interaction is basically determined by the electron temperature in the laser-plasma interaction region. The results shown in Fig.5.3.1 and Table 5.3.1 have been obtained assuming an electron temperature of 2 keV which, consistently with the limits of the 1-D simulation discussed above, is lower than the 6 keV suggested by MEDUSA. However, although suggested by 2-D numerical simulation, the particular choice of 2 keV was somewhat arbitrary.

Nevertheless, the conclusions obtained here are basically independent from the particular choice of the electron temperature. In fact, according to Eq.3.3.1, for higher values of the electron temperature the E-field strength coefficient becomes smaller and the effective temperature introduced in Eq.3.3.3 approaches the thermal value. In particular, in the case of the 1-D model, the predicted electron temperature is high enough that non-linear effects can be neglected. At this stage of the investigation, the balance between thermal and non-linear effects in the absorption process is still undefined. However, an effective temperature in the laser-plasma interaction region of typically 5 keV allows the whole set of absorption measurements presented here to be fully consistent and satisfactorily modelled.
5.4 - A Study of Plasma Self-Emission

A more comprehensive characterization of the plasma produced by irradiation of solid targets with the 12 ps, 268 nm SPRITE laser beam requires a study of the equilibrium conditions as discussed in detail in Sect.2.2 and Sect.3.4. According to the simulation and the conclusions of the previous section, electron temperatures of the order of several keV can be expected with consequently strong emission of X-ray radiation. An extensive investigation was carried out on X-ray and X-UV plasma self-emission using time-resolved spectroscopy, and the results were compared with the predictions of atomic physics numerical models.

**Experimental set-up**

A schematic view of the arrangements of the main diagnostics employed in this study is shown in Fig.5.4.1. The laser focusing configuration was analogous to that shown in Fig.5.1.1. A flat field grazing incidence, 2400 l/mm grating (Nakano et al., 1984) was employed in order to analyse plasma self-emission in the spectral range from 20 to 80 Å.

![Time Resolving X-ray Spectrometer](image)

**Fig.5.4.1.** Schematic arrangement of X-ray diagnostics for time integrated and time resolved spectroscopic measurements. The optics used to focus the laser beam on target was the same as described in Fig.5.1.1 and Sect.5.1.

This grating has a variable groove spacing which allows the dispersed radiation to be focused onto a plane rather than a curved surface as in standard gratings. Such a feature makes this grating particularly suitable for use with flat detectors such as
film plates or streak-camera photocathodes. A hard X-ray filter was also used with the grating in order to cut off high energy photons (>1keV) for measurements in the range above ≈15 Å. In fact this energetic radiation, if allowed on the grating, would contribute to the first order spectrum with higher diffraction orders which would strongly reduce the signal-to-noise ratio.

The high energy cut-off filter was developed by employing the properties of X-ray reflection at grazing incidence angles. A double reflection mirror was implemented using two highly parallel, highly polished, Ni-coated silica substrates. The unit was placed on the path of the radiation before the grating as shown in Fig.5.5.2. The L-shell absorption edge of Ni atoms, at approximately 860 eV, defines the low energy cut-off edge of the filter. Below this value the reflectivity of the Ni-coated surface is very close to unity (≈ 95 %) provided the incidence angle is ≤ 50 mrad.

![Fig.5.4.2. Schematic diagram of the experimental arrangement of the flat-field XUV grating showing the double reflection filter used to cut off high energy (>1 keV) photons. The X-ray streak-camera was replaced with X-UV film plates for time integrated measurements.](image)

Time resolved measurements were carried out by coupling the spectrometer to an X-ray streak-camera. The photocathode of the streak-camera consisted of a 0.1 µm CH foil coated with a layer of photo-sensitive material. CsI was typically used for low temporal resolution (>10 ps) measurements while, when a temporal resolution better than 10 ps was required, KBr was used. Although less sensitive than CsI (Chekovin, 1990), KBr photocathodes allow a better temporal resolution as the secondary electrons emitted have a narrower energy distribution (Henke et al., 1979). For time integrated measurements, spectra were recorded using an X-UV sensitive film (Kodak 101) placed at the output plane of the spectrometer.

Finally, time integrated images in the X-ray spectral region were recorded using an active pin-hole camera fitted with a four pin-hole array as described in Sec.4.4. Each imaging channel was set to be sensitive to different spectral regions by using different X-ray filters. One channel was filtered with a 20 µm thick Be filter, while the other three channels were additionally filtered with 3 µm Al, 8 µm Cu, and 3 µm...
Al+8 µm Cu respectively. The transmittivity of a 25 µm Be, 3 µm Al and 8 µm Cu foil is plotted in Fig.5.4.3 as a function of the X-ray photon energy in the range from 100 eV to 100 keV, calculated using available (Henke et al., 1985) mass absorption coefficients.

![Transmittivity vs Energy](image)

**Fig.5.4.3.** X-ray transmittivity of a 3 µm thick Al foil, a 25 µm thick Be foil and a 8 µm thick Cu foil calculated using experimentally measured mass absorption coefficients.

According to Fig.5.4.3, the imaging channel filtered with the Be foil is therefore sensitive to photon energies above ≈1.5 keV. The addition of the Al filter in the second channel further shifts this lower bound to approximately 2.5 keV. In the case of X-ray emission from Al plasmas, the Al filter also cuts off most of the line emission from highly ionized Al ions, as already shown in Fig.4.4.4. The second and the third imaging channels are mainly characterized by the transmittivity of the Cu foil. In this case the presence of the L-shell absorption at 933 eV selects photon energies between ≈6 keV and 9 keV and greater than ≈15 keV.

The effect of the different filtering conditions on the image of the X-ray plasma sources is shown in Fig.5.4.4. The plasma was produced irradiating a layered target (see Sect.5.5) consisting of a solid plastic (Mylar) target coated with a 0.1 µm thick Al layer and overcoated with a 0.1 µm thick plastic (CH) layer. The laser beam was focused in a ≈25 µm diameter focal spot, at an intensity of 3.6×10^{16} W/cm². The dynamics of the laser-target interaction in this special configuration will be discussed in detail in Sect.5.5. The plots of Fig.5.4.4 were obtained from one-dimensional line-outs taken along the diameter of each of the four disk-shaped X-ray images. The horizontal axis is labelled in terms of the distance on the target plane. It is instructive to compare the size of the focal spot given in Sect.5.1, measured using the obscuration technique, Sect.5.1, with the width of the profiles of Fig.5.4.4.
Fig. 5.4.4. Lineout of four differently filtered X-ray images obtained on a single event using a pin-hole camera fitted with a 4-pin-hole array. The plasma was produced by irradiating a solid plastic (Mylar) target coated with a 0.1 µm thick Al layer, overcoated with a 0.1 µm thick plastic (CH) layer. The laser intensity was $3.6 \times 10^{16}$ W/cm$^2$ in a $\approx 20$ µm diameter focal spot.

The FWHM of the profile relative to the Be filtered channel is approximately $\approx 90$ µm while the Be+Al and the Be+Cu both give a FWHM of 60 µm. Finally the Be+Al+Cu channel gives a FWHM of 30 µm. According to these results the extent of the X-ray emitting region, in the case of the Be+Al+Cu channel, is limited to the $\approx 25$ µm diameter focal spot of the laser. On the other hand, according to Fig.5.4.3, only very energetic X-rays contribute to this image. These X-rays can only be produced by continuum emission, bound-bound processes being energetically limited to the maximum transition energy of the mostly ionized atom present. In the particular case considered here, the maximum bound-bound transition energy in H-like Al ions, is given by the ionization energy of the H-like Al ion which, according to the plot of Fig.4.4.4, is $\approx 2.3$ keV, that is well below the lower energy bound of $\approx 6$ keV set by the transmittivity of the Cu filter.

Continuum emission, and in particular that arising from free-free transitions (bremsstrahlung), is directly related to the electrons heated by the laser e.m.wave and, therefore, provides a good indication of the actual extent of the laser intensity distribution on target. On the other hand, an accurate evaluation of this distribution also requires that lateral heat flow is taken into account. This aspect has already been discussed in Sect.5.3 and, according to the conclusions drawn therewith, little contribution can be expected from lateral heat transport to the transverse size of the hot plasma in the region close to the critical density where most of the absorption takes place and where high energy photons are therefore produced.
It is interesting to note that the FWHM of the profile taken with the 8 µm Cu filter only is 60 µm, i.e. two times larger than the 30 µm considered above where a 3 µm Al filter was added to the Cu filter. Moreover, a similar effect can also be observed comparing the result obtained with the Be filter only, which gives a FWHM of ≈90µm which is three times larger than the FWHM obtained when the Al filter is added to the Be filter. These results indicate that the region immediately outside the laser focal spot exhibits a strong X-ray emissivity in the spectral range between ≈1.5 and ≈3 keV, where the Al filter mainly contributes to the attenuation, due to the K-shell absorption edge at 1560 eV. As already pointed out above, in the particular case of the target considered here, the presence of the Al layer will definitely give rise to strong line emission arising from bound-bound transitions in He-like and H-like Al ions. Whether this region is directly heated by the laser energy present in the wings of the laser intensity distribution, or is heated due to lateral electron heat conduction is still an open issue. However, it is likely that both effects contribute, in our experimental conditions, to distribute the laser energy over the observed region.

**Time integrated spectra of low Z highly ionized ions.**

A preliminary estimate of the degree of ionization of the plasmas produced by irradiation with the 268 nm, 12 ps laser pulse can be obtained from spectroscopic investigation of the X-ray radiation emitted by the plasma. Fig.5.4.5 shows a lineout of the spectrum obtained by irradiating a solid plastic (Mylar) target at an intensity of ≈2 × 10^{14} W/cm^2 in a 200 µm diameter focal spot.

![Graph showing X-ray spectrum](image)

**Fig.5.4.5.** Time integrated spectrum of the X-ray radiation from the plasma produced by irradiation of a solid plastic (Mylar) target with a 12 ps, 268 nm laser pulse at an intensity of ≈2×10^{14} W/cm^2 in a 200 µm diameter focal spot.
The spectral range of the spectrometer was set to detect line emission from hydrogenic-like and He-like C ions. The film density was converted into intensity using the calibration data available (Schwanda & Eidmann, 1992) for the Kodak 101 film used in these measurements. The intensity was also normalized to the maximum intensity in the spectrum. Spectral calibration was carried out from comparison with the results of the simulation of the atomic physics using the numerical code RATION.

The most striking feature of this spectrum is the predominance of hydrogen-like carbon emission over the He-like emission. In particular, the Lyα line, i.e. the 1s-2p transition of hydrogen-like Carbon, is 4.1 times more intense than the line relative to the 1s²-1s2p transition of He-like Carbon (C Heλ). These circumstances provide a first indication on the average plasma emission temperature when one takes into account that the ionization energy of C, He-like ions is 392.1 eV. It is instructive to compare these results with the analogous ones obtained from interaction with a material characterized by a different atomic number. The spectrum shown in Fig.5.4.6 was obtained at the same irradiance as that of Fig.5.4.5 but with a solid Al target. In this case a multi-configuration Dirac-Fock atomic physics code (Grant et al., 1980) was also used in the identification of the spectral lines, in order to calculate the oscillator strength of the transitions relative to Be-like ions, not included in the RATION package.

![Fig.5.4.6.](image)

**Fig.5.4.6.** Time integrated spectrum of the X-UV radiation from plasmas produced by irradiation of solid Al with a 12 ps, 268 nm laser pulse at an average laser intensity of ≈2×10¹⁴ W/cm² in a 200 µm diameter focal spot.

According to the results of the hydrodynamic simulation performed using MEDUSA for the interaction conditions of Fig.5.4.6, the electron temperature at the critical electron density for the 268 nm laser light, i.e. at \( n_e \approx 1.5 \times 10^{22} \text{ cm}^{-3} \), is
expected to be approximately 400 eV. The synthetic spectrum generated by RATION for these conditions, although not inclusive of the Be-like transitions, is shown in Fig.5.4.7 for the same spectral range of Fig.5.4.6.

![Figure 5.4.7](image_url)

**Fig. 5.4.7.** Synthetic spectrum of the X-ray emission in the spectral range of Fig.5.4.6 for a 1 µm thick Al plasma in an optically thick regime, at an electron density of $1.5 \times 10^{22}$ cm$^{-3}$.

The intensity of the hydrogenic, He-like and Li-like lines was calculated assuming a 1 µm plasma size in an optically thick regime, where optical depth effects are approximated by escape factors (Mihalas, 1978). A comparison between the experimental spectrum and the calculated one suggests that the electron temperature predicted by the hydrodynamic simulation is consistent with the experimental data. However, this comparison takes no account of the space/time integration effects (Desselberger et al., 1992) which characterize the experimental data of Fig.5.4.6. This aspect will be discussed below.

**Space integration effects on XUV spectra**

In order to estimate the importance of space integration on the spectra like that of Fig.5.4.6, the distribution of the ion species of interest, namely Hydrogenic, He-like, Li-like and Be-like, was also calculated by the hydrodynamic simulation using an Average Atom model, with the same interaction conditions considered above. **Fig.5.4.8** shows this distribution as a function of the distance from the target surface at the peak of the laser pulse. The calculated electron temperature is also plotted on the same graph. The main feature of this graph is the abundance of He-like ions over the whole plasma extent. In contrast, hydrogenic, Li-like and
Be-like ions are well localized in space and are characterized by values of the local electron temperature which can be significantly different from each other. In fact, according to this calculation, the electron temperature at the maximum of the hydrogenic ion density is approximately 450 eV whereas the Li-like ion density peaks at an electron temperature of 180 eV. A further contribution to spatial integration effects may also come from any lateral heat flow that may occur during the laser-target interaction as already discussed above and in Sect.5.3.

![Graph](image_url)

**Fig.5.4.8.** Ion density distribution for hydrogenic, He-like, Li-like and Be-like ions calculated at the peak of the pulse, for irradiation of a solid Al target with a 268 nm, 12 ps laser pulse at an intensity of \(\approx 2 \times 10^{14}\) W/cm\(^2\). The electron temperature is also plotted (dashed curve).

In summary, these results indicate that, in highly transient plasmas as those investigated here, the comparison between line intensities relative to different ionization stages, as a diagnostic for electron temperature measurements, must take into account spatial integration effects, unless spectroscopic measurements with spatial resolution are carried out. The dynamics of heat transport processes and their consequences on radiation emission processes will be discussed in the next section where measurements of mass ablation rates in a highly transient regime are presented.
5.5 - Transient Heat Propagation

The production of hot high density plasmas is of great significance to both basic atomic physics studies and X-ray laser research (Murnane et al., 1991, Rosen, 1990, Smith et al., 1990). These plasma conditions can be achieved when a prepulse-free picosecond laser pulse is focused onto solid targets. Since these values of electron density are well above the critical density, where laser light cannot propagate, thermal conduction plays a key role in the transport of energy from the laser absorption region to the high density plasma region. A better understanding of thermal transport mechanisms is therefore vital in the study of these plasmas. In the experiment reported here, mass ablation rate measurements were performed in a picosecond regime. X-UV time-resolved spectroscopy with picosecond temporal resolution was applied for the first time to burn-through measurements. A detailed characterization of heat propagation in hot high dense plasmas generated by picosecond laser pulses is presented.

Heat transport in laser-produced plasmas

The modelling of heat propagation in the simulation of laser-plasma interaction experiments is presently being performed by mainly using the classical flux-limited Spitzer-Harm thermal conductivity (Spitzer & Harm, 1953). On the other hand, detailed 2D Fokker-Planck(FP) codes have already been successfully applied (Rickard et al., 1989) to the interaction of high intensity laser light with hot dense plasmas. Nevertheless, the implementation of laser interaction codes with FP routines, to describe thermal conductivity, is presently not feasible due to the large amount of computer-time needed. A comparison of the results of these two different approaches, when applied to simple physical systems, can however provide information on the constraints as well as the range of validity of the classical flux limited description. Moreover, the best value of the phenomenological parameter \( f_L \), the so called flux limiter, can be inferred for the particular plasma condition under investigation.

Intense experimental activity has been devoted to investigating thermal electron transport in laser produced plasmas in both planar and spherical geometry (Dahmani et al., 1991). Since a direct measurement of thermal electron flux is not
possible, physical quantities depending on the thermal transport are usually considered and measured (Krueer, 1988). The mass ablation rate, \( m \, (\text{g cm}^{-2} \, \text{s}^{-1}) \) is usually taken into account (Key et al., 1983) as it is relatively easy to evaluate this quantity from burn-through studies in multi-layered targets or from ion velocity measurements using Faraday detectors. The propagation of an isotherm of the heat front can be revealed by studying X-ray emission whose spectral features are characteristic of each layer of the target.

Most of the works published so far have been concerned with long (\(<\text{ns}\)) laser pulses where quasi-steady-state heat propagation can be assumed. Temporal dependence of the mass ablation rate has also been studied in UV-laser-irradiated spherical targets (Jaanimagi et al., 1986). In that investigation the authors pointed out that time-dependent effects occurring during the laser pulse can significantly affect thermal transport mechanisms. An experimental investigation on heat transport in a transient regime was carried out using the 12 ps laser pulse described in the previous sections and the results are presented and discussed below.

**Time resolved X-UV spectroscopic measurements**

The 12 ps, UV laser pulse was focused onto a three-layer planar target at irradiances of up to \( 3 \times 10^{16} \, \text{W/cm}^2 \), that is, in a regime similar to that at which the absorption measurements presented in Sect.5.3 were performed. The target configuration is schematically shown in Fig.5.5.1. It consisted of a \( 10 \, \mu\text{m} \) thick Mylar substrate coated with 2000 Å Aluminium and overcoated with a thin (1000-2000 Å) low Z plastic ablator (CH).

![Fig.5.5.1. Schematic configuration of the multi-layered target used for mass ablation measurements in a transient picosecond regime. The 12 ps, 268 nm laser pulse was incident on the target at an intensity ranging from \( 10^{15} \) to \( 5 \times 10^{16} \, \text{W/cm}^2 \).](image)

Mass ablation rates were inferred from the history of X-UV line emission generated by the signature layers as the heat front propagated in the target. The
primary diagnostic was the time-resolving XUV spectrometer already described in Sect. 5.4. A grazing incidence flat field grating was coupled to a newly developed X-ray streak-camera. As already described in Sect. 5.4, a Ni-coated reflection filter (see Fig. 5.4.2) was placed before the grating in order to cut-off high energy photons (>1keV). The spectral range of the spectrometer was chosen in order to detect characteristic line emission from each one of the three layers. Line emission from H-like Carbon (C-Ly\( \alpha \)) from the C-H layer, Li-like Aluminium (Li 1s\(^2\)2p\(^-\)1s\(^2\)4d) from the middle Al layer and both H-like Carbon (C-Ly\( \alpha \)) and H-like Oxygen (O-Ly\( \alpha \)) from the Mylar substrate were detected in the same spectra. The streak-camera employed to temporally resolve the spectrum was capable of a sweep-speed as high as 3.5 ps/mm at the output phosphor window and was fitted with a KBr photocathode which resulted in a temporal resolution of \( \approx 1 \) ps.

The temporal evolution of the above specified Oxygen, Carbon and Aluminium spectroscopic lines obtained at an irradiance of \( 3 \times 10^{16} \) W/cm\(^2\) is shown in Fig. 5.5.2. The first peak of C-Ly\( \alpha \) emission arises from the first C-H layer. As the temperature in this region rises, Carbon atoms soon become fully stripped and emission intensity falls. Then the heat front propagates in the Al layer giving rise to Li-like Al emission. Finally C-Ly\( \alpha \) emission takes place again when the inner Mylar substrate is heated.

![Fig. 5.5.2.](image)

**Fig. 5.5.2.** Temporal evolution of X-UV line emission from laser irradiation of a CH-Al-Mylar target at an irradiance of \( 3 \times 10^{16} \) W/cm\(^2\). Line emission from H-like Carbon (C-Ly\( \alpha \)), from the C-H layer, Li-like Aluminium (Li 1s\(^2\)2p\(^-\)1s\(^2\)4d) from the middle Al layer and both H-like Carbon (C-Ly\( \alpha \)) and H-like Oxygen (O-Ly\( \alpha \)) were detected simultaneously.

Due to the nominal target composition and configuration, emission from Oxygen ions should resemble the second peak of carbon emission. However, though it lasts longer than the Al emission, it starts almost simultaneously with such emission from
the middle layer. This can be explained by the presence of Oxygen impurities in the Aluminium layer due to the presence of Al oxide (Al₂O₃). In fact strong Oxygen emission was systematically detected when solid "pure" Aluminium targets were irradiated.

Assuming that the history of the X-UV emission shown in Fig.5.5.3 mainly depends upon the heat propagation and that atomic processes are fast compared to the typical time-scale of these heating mechanisms, then one can immediately determine the mass ablation rate from the plot of Fig.5.5.2. Considering the density and the thickness of the ablated Al layer and the peak-to-peak distance in the history of C-Lyα which is 22 ps ± 2ps one obtains a mass ablation rate of \[ m = 2.4 \times 10^6 \text{ g cm}^{-2} \text{ s}^{-1} \pm 10\% \]. A similar procedure was used to determine the mass ablation rate relative to a lower incident irradiance of \( 4.3 \times 10^{15} \text{ W/cm}^2 \). In this case we found \[ m = 3.4 \times 10^6 \text{ g cm}^{-2} \text{ s}^{-1} \pm 10\% \]. A detailed analysis of the experimental results as well as a detailed discussion of all the processes which can affect this simple assumption is given below.

**Transient effects in X-UV emission processes**

The main feature of Fig.5.5.2 is the temporal evolution of C-Lyα line which is characterized by two maxima separated by 22 ps ± 2ps. The first peak arises from the interaction of the laser energy with the first CH layer. When the ablation front reaches the plastic (Mylar) substrate, after propagation in the Aluminium layer, C-Lyα emission takes place, again resulting in the second peak of emission.

We observe that in order for this radiation to be a good diagnostic for the measurement of heat propagation atomic physics processes must be fast compared to the hydrodynamic ones. In the case that this condition is not fulfilled, a time-dependent atomic physics code would be necessary. In the experimental investigation on mass ablation rate published so far where nanosecond laser pulses were considered, a similar method as that described so far was used. The ablation rate was inferred from the temporal evolution of H-like Al or H-like Si lines. As already pointed out in Sec.3.4, according to the plots of Fig.3.4.1-2, at the electron densities and temperatures of laser produced pulses in the nanosecond regime, ionization times of several hundreds of picosecond can be expected. Therefore the accuracy of those mass ablation rate measurements from time resolved X-ray spectroscopy of highly ionized, medium Z plasmas, is likely to be affected by time dependent effects. Regardless of the particular temporal regime under investigation, time dependent effects on mass ablation rate measurements carried out using X-ray
emission as a diagnostic should always be carefully examined. On the other hand, one may regard transient atomic physics effects as giving rise to constraints on the particular mechanism of radiation emission to be selected, as suitable for the particular temporal scale under investigation.

Concerning the experiment described here, according to the general criteria outlined in Sect.3.4 and considering that we are dealing with emission from low Z elements (Carbon and Oxygen) or Li-like Al, we expect the condition for a steady-state model to be fulfilled. For a more quantitative analysis the relaxation time for ionization from He-like to H-like Carbon, He-like to H-like Oxygen, Be-like to Li-like Aluminium and He-like to H-like Aluminium is reported in Table 5.5.I as calculated according to Eq.3.4.7

<table>
<thead>
<tr>
<th>Te = 400-800 eV</th>
<th>C^{4+}-C^{5+}</th>
<th>O^{6+}-O^{7+}</th>
<th>Al^{9+}-Al^{10+}</th>
<th>Al^{11+}-Al^{12+}</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \tau ) (ps)@ne=1.5E22 cm(^{-3})</td>
<td>0.4 - 0.2</td>
<td>2.7 - 0.9</td>
<td>0.4 - 0.2</td>
<td>49.2 - 22.7</td>
</tr>
<tr>
<td>( \tau ) (ps)@ne=1.5E23 cm(^{-3})</td>
<td>0.4-0.02</td>
<td>0.26-0.07</td>
<td>0.04 - 0.02</td>
<td>4.7 - 2.2</td>
</tr>
</tbody>
</table>

Table 5.5.I. Relaxation time (See Eq.3.4.7) for ionization from He-like to H-like Carbon, He-like to H-like Oxygen, Be-like to Li-like Aluminium and He-like to H-like Aluminium at two values of the electron density relevant to the experiment described here. The left-hand (right-hand) value refers to an electron temperature of 400 eV (800 eV).

According to these results, all the relaxation times of interest in our experimental condition are a fraction of a picosecond, with the exception of the Oxygen at the lower value of electron density considered in the table. The behaviour of the ionization dynamics of Oxygen ions will be discussed later in this section as it can contribute to explain the apparently anomalous temporal evolution of O-Ly\( \alpha \) emission. Another peculiarity of the results presented in Table 5.5.I is that the relaxation time to H-like Aluminium ions is typically two orders of magnitude larger than the other values, as it ranges from a few picoseconds to several tens of picoseconds. This is mainly due to the large increase in the ionization energy of K-shell electrons. This fact has an important consequence on the ionization dynamics in the transient phase. In fact, as the Al layer is initially heated, its temperature rises, ionization takes place and in less than a ps, He-like ions are generated. The large ionization energy of He-like ions however, will slow down further ionization making the Li-like/He-like system a temporary equilibrium condition. As time passes and the electron temperature increases further, the equilibrium will eventually shift towards a He-like/H-like system. This process makes Li-like and He-like ions particularly suitable as steady state sources of radiation during the first picoseconds of the plasma start-up phase.
Finally, we consider the time-scale of radiative spontaneous emission for the three transitions considered above, given by the Einstein coefficient of Eq. 3.4.6. From the available oscillator strengths (Martin & Wiese, 1983) we obtain decay times of 1.6 ps, 0.5 ps and 3 ps for C-Lyα, O-Lyα and Li 1s²2p-1s²4d line emission respectively. According to this analysis we can conclude that, although transient effects in radiation emission can give a minor contribution to the observed history of the line emission the main features of the X-ray emission processes taken into account in this study are expected to be mainly related to the hydrodynamic evolution of the plasma.

**Computer modelling of the mass-ablation rate**

The hydro-code MEDUSA was used to simulate the interaction of high intensity 268 nm laser radiation with planar targets. The general features of the physics included in the code have already been given in Sect.2.1, while a general overview of the numerical simulation of short-pulse experiments has been presented in Sect.3.1. The results reported in this section are of particular relevance to the mass ablation rate measurements. A 12 ps (FWHM) Gaussian laser pulse was incident onto a three-layer planar target at irradiances up to $3 \times 10^{16}$ W/cm². The target consisted of a 10 μm thick Mylar substrate coated with 2000 Å Aluminium and overcoated with a thin (1000 Å) plastic layer (CH).

Laser energy was absorbed via inverse bremsstrahlung with a fraction of the remaining energy being damped at the critical surface to simulate anomalous absorption. The resonance absorption coefficient in the code was set according to the results of the absorption measurements described in Sect.5.2 and Sect.5.3. Although many factors cooperate to affect the efficiency of resonance absorption on a shot-to-shot basis such as target roughness or beam quality for example, simulation shows that the effect of variation of resonance absorption on the ablation rate can be ignored if variation of up to approximately 20% are introduced into the absorption coefficient. As already discussed in Sect.2.1, the harmonic mean between the classical heat flux, and the free streaming value, is used by the code (See Eq.2.1.4) to limit the heat flux. A total absorption of approximately 18% at the peak of the pulse and 21% off peak is accounted for by the code in the high incident laser irradiation regime. This value is consistent with the plot of Fig.5.3.1 which, at a slightly higher incident irradiance of $5 \times 10^{16}$ W/cm², gives a total absorption of approximately 20%. The propagation, in the Lagrangian reference frame, of an isotherm of the heat front at a given temperature, is used to evaluate the
mass-ablation rate from the results of the simulation. Electron temperatures between 200 eV and 500 eV were considered in this analysis, consistent with the range of photon energies measured experimentally as described in detail below. The mass ablation rate was determined according to its definition

\[
\dot{m} = \rho \frac{\Delta r}{\Delta t}
\]  

(5.5.1)

where \( \rho \) is the solid density of the ablated material, \( \Delta r \) is the burn-through depth, \( \Delta t \) is the time needed by the burn-through to take place. \textbf{Fig.5.5.3} shows the temporal evolution of electron temperature in the C-H and Mylar layers for an incident laser intensity of \( 3 \times 10^{16} \text{ W/cm}^2 \)

![Electron Temperature vs Time](image)

\textbf{Fig.5.5.3.} Temporal evolution of electron temperature in the C-H and Mylar layers as calculated by MEDUSA for an incident laser intensity of \( 3 \times 10^{16} \text{ W/cm}^2 \) and with a resonance absorption factor of 10%.

The temperature in the C-H layer rises to several hundred eV in approximately 1 ps. After propagation in the middle layer (Al), the heat front reaches the Mylar substrate. The electron temperature in the Mylar substrate at the first 10 cells allocated to this layer in the Lagrangian reference system is plotted in the graph of Fig.5.5.3. A minimum delay \( \Delta t \approx 9 \text{ps} \) is predicted by the code between the time at which the CH layer is heated at 400 eV and the time at which the first cell of the Mylar substrate reaches the same temperature. We observe that the dependence of such a delay upon the particular value of temperature chosen is weak in the range taken into account. This is consistent with the steepness of the heat front predicted by the classical theory of heat propagation. From the plot of Fig.5.5.3 one can derive mass ablation rates by measuring the time interval, \( \Delta t \) for a given set of input
parameters including laser intensity, heat flux limiter and resonance absorption factor. The effect of resonance absorption on such results has been examined in the simulation by varying the amount of laser energy dumped at the critical layer. Fig. 5.5.4 shows an output analogous to the one of Fig. 5.5.3 but with the anomalous absorption factor reduced from 10% to 5%.

![Temporal evolution of electron temperature in the C-H and Mylar layers as calculated by MEDUSA for the same irradiance as that of Fig. 5.5.3 and with a resonance absorption factor of 5%.

Although in this case a lower electron temperature is produced in the inner Mylar layer, the time needed by the heat front (at 400 eV) to reach this layer is $\Delta t \approx 10$ ps which differs from the previous value relative to Fig. 5.5.3 by only 1 ps. On the other hand, from an experimental viewpoint, the temporal resolution of our measurements is approximately 1 ps. Therefore, fluctuations of the absorbed laser energy due to resonance absorption, within the above specified range, would produce variations of mass ablation rate within the experimental error.

Finally, the dependence of mass ablation rates on the absorbed laser irradiance for different values of the flux limiter was investigated. Fig. 5.5.5 shows the mass ablation rate as a function of absorbed laser intensity for two values of the flux limiter, i.e. $f_L = 0.1$ and $f_L = 0.04$. The mass ablation rate was calculated for several values of the incident intensity as indicated by the data points shown in the graph. The same values of the incident intensity were chosen for the two values of the flux-limiter. The solid line was obtained from a simple fit to the numerical data. Notice that the ablation rate has been plotted as a function of the absorbed intensity. This plot clearly shows that the heat flux limiter plays a very important role in these interaction conditions. In fact, according to the simulation, the ablation rate increases by 25% when the flux limiter changes from 0.04 to 0.1.
Fig. 5.5.5. The mass ablation rate as a function of the absorbed laser intensity, calculated by MEDUSA for two different values of the flux limiter. A layered (CH-Al-Mylar) target was irradiated with a 12 ps gaussian pulse and the average ablation rate of the middle Al layer is reported here. The result of the experimental measurement is also plotted in the same graph for comparison.

As a consequence of this increased flow of thermal energy toward the higher density region, the subcritical plasma, where absorption occurs, becomes cooler and more collisional and the absorption increases by typically 20%, for the same change of heat flux. We stress here that this dependence upon the flux-limiter is a strong indication of the fact that the classical theory of heat conductivity cannot properly describe this regime of interaction. In fact, as explained in detail in Sect.2.1, the flux limiter starts playing a role when the temperature gradients are so large that the classical theory would lead to non-physical results. An upper limit to the heat flux is therefore imposed using this phenomenological parameter.

Conclusions

As shown by Fig.5.5.5, the effect of this limitation increases as the laser intensity increases and larger temperature gradients are established in the plasma. The comparison of the simulation with the experimental results clearly shows this effect. The measured values of mass ablation rate given above for the incident intensities of $3 \times 10^{16} \text{ W/cm}^2$ and $4.3 \times 10^{15} \text{ W/cm}^2$ are also plotted in Fig.5.5.4. The incident intensity was converted into absorbed intensity according to the results presented in Fig.5.2.3, which give the absorption as a function of the incident intensity measured in similar conditions to that under investigation here. The error bars in the ablation rate were assumed to be 20% instead of the 10% resulting from the temporal resolution of the data. In fact, the conclusions on the temporal response
of the X-ray lines reached above suggest that transient effects in the emission processes could be another source of error.

The mass absorption coefficient measured in the lower intensity case is consistent with the calculated one for a flux limiter of 0.04. This value, though not very different, is smaller than that inferred by the 2D Fokker-Plank simulation performed in a similar intensity regime, as already discussed in detail in Sect. 5.3. In that study it was found that the heat flux in the subcritical region was consistent with a classical value using a flux limiter of 0.1. However, we cannot exclude that other sources of error, not accounted for in this analysis, could make this experimental result consistent with the 2D FP simulation.

In contrast with the consistency between the calculated and the measured ablation rate in the lower intensity regime, a large discrepancy is found in the higher intensity regime, where the measured ablation rate is less than half of the calculated one. This would indicate a strong reduction of the heat flux compared with the lower intensity case. It is clear that this discrepancy cannot be simply explained in terms of failure of the classical heat transport theory assumed in the calculation as this conclusion would lead to an unacceptable value of the heat flux limiter, orders of magnitude smaller than the free streaming value. Several effects can be invoked in order to explain this result. On the other hand it has been suggested (Riley et al., 1992) that strong self-generated magnetic fields arising from the $\nabla n_e \times \nabla T_e$ mechanism (Haines, 1986), could lead to substantial inhibition of the thermal transport if the Larmor radius of the streaming electrons is small compared to their mean–free path. It is clear however that a complete modelling of the interaction in this high intensity regime would require a self-consistent treatment of non–local heat transport and magnetic field generation in a two dimensional geometry.

Nevertheless, as already pointed out in Sect. 5.3, lateral heat flux can play an important role in the subcritical region, where laser-plasma interaction occurs. This aspect has already been extensively discussed in the analysis of the absorption measurements and the conclusions reached there are fully consistent with the results being discussed here. In fact, it was pointed out that in this interaction regime, the conditions may be fulfilled so that free streaming lateral heat flow could occur in the subcritical plasma region. Since this is a purely 2D effect, the 1D calculation presented here would obviously fail to account for it. From the point of view of the experimental results, any lateral heat flow would result in less thermal energy being transported into the inner layers of the target, thus effectively reducing the mass ablation rate.
References


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List of Author’s Publications Related to the Ph.D Project


