SHOCK WAVES GENERATED BY INTENSE FEMTOSECOND LASERS

By

Tommy Ao

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Department of Physics and Astronomy

The University of British Columbia

6224 Agricultural Road

Vancouver, B.C., Canada

V6T 1Z1

Date:

SEPT. 7, 2000

Abstract

The advent of intense femtosecond lasers has created the exciting possibility of accessing regimes of extreme high pressure using a relatively small laser system. This stems from the lack of significant hydrodynamic expansion during the process of laser deposition in a solid via skin-depth absorption, which leads to extremely high energy densities in the irradiated sample. After the short-pulse laser energy has been absorbed, the laser-heated material begins to be released which drives a shock wave into the sample. However, unlike previous long-pulse laser driven shock waves, the shock wave driven by a intense short-pulse laser rapidly decays as it propagates through the sample. Before adopting such a shock wave as a new approach in the study of high density plasmas, its unique characteristics must be understood.

A one-dimensional hydrodynamic code which is coupled to an electromagnetic wave solver is used to elucidate the basic properties of shock waves generated by intense femtosecond lasers. Using a unique experimental scheme, the electrical conductivity of silicon in the dense, plasma state can also be studied. Calculations were performed in which a shock wave was driven into a silicon sample by a pump laser with a wavelength of 400 nm, pulse length of 120 fs (FWHM) and irradiances ranging from $10^{14} - 10^{15} \text{ W/cm}^2$, while rear-side optical measurements were made by a 800 nm, 120 fs probe laser.

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Chapter 1

Introduction

1.1 Strongly Coupled Plasmas

Within the field of plasma physics the area of strongly coupled plasmas (SCP) remains the least understood. The study of these plasmas has been motivated by research in geophysics, astrophysics, x-ray sources, and nuclear explosions, to name a few. Yet, the main driving force behind the study of SCP continues to be the goal of achieving inertial confinement fusion. In this regime of high density and modest temperature, the plasma cannot be described solely by the statistical behaviour of a screened Coulomb system as in the case of a rarefied plasma; effects due to strong particle correlations must be addressed. Simply stated, when the potential energy due to the Coulomb interaction between the ions is greater than kinetic energy of the ions, the plasma is considered strongly coupled.

In order to create a strongly coupled plasma, great pressures must be achieved. The methods employed are usually divided under the headings of static pressure-loading and dynamic pressure-loading. The most advanced static techniques involve the use of diamond anvil cells to slowly compress a solid. However, the highest pressure achievable in this manner (< 2 Mbar; $1 \text{ TPa} = 10 \text{ Mbar}$) is ultimately limited by the finite yield-strength of the diamond anvils [1]. Thus, in order to reach higher pressures dynamic pressure-loading methods must be used. The basic principle in dynamic techniques involves the generation of a rapid compression (shock) wave in the material to be studied. In the early days of shock physics, explosive systems such as shock tubes were employed

to achieve pressures of a few hundred kbars [2, 3, 4]. This method was later refined into the use of explosive-driven flyer-plates allowing precise measurements ($< \pm 2\%$ pressure uncertainty) to be made at pressures near 10 Mbar [5, 6]. The current state-of-the-art precision measurements are obtained using two-stage light-gas guns [7]. Even higher pressures can be achieved via nuclear explosions (hundreds to thousands of Mbars) [8]. However this method is quite expensive and unavailable to most researchers, not to mention the obvious political and environmental ramifications associated with it.

1.2 Laser-Driven Shock Waves

With the advent of high-power pulsed lasers in the early 1970's, a new avenue for achieving ultra-high pressures became available (10-100 Mbar) [9, 10]. Much of this early research involved heating and compressing small fuel pellets in inertial confinement fusion schemes [11, 12, 13]. In fact, current advances in high-power pulsed laser technology remain mainly due to the work on achieving inertial confinement fusion. However, other uses for these lasers soon became evident, and over the past two decades an array of research has been carried out [14, 15, 16]. In addition to the advantage of reaching pressures unattainable with the other methods, the laser-driven shock technique has the added advantage of being able to easily control the pressure by simply varying the laser irradiance on the target. This technique is intrinsically different from those previously described in several key aspects. Unlike the other methods where the shock is produced by an impact or explosion, it is formed through the laser ablation (vaporization) process. This process is discussed below, but essentially the shock arises out of conservation of momentum. Experiments of this nature are inherently spatially microscopic and short-lived (less than a few nanoseconds), thus requiring diagnostics with high spatial and temporal resolutions.

Experiments involving high-power pulsed lasers to create a strongly coupled plasma

can be separated according to the pulse-width of the laser used. Prior to the mid 1980's, all high-power pulsed laser research employed lasers with pulse-widths ranging from tens or hundreds of picoseconds up to nanoseconds, and are considered "long"-pulse lasers. In these long-pulse laser experiments, as the target is irradiated by the laser pulse, the ablated surface expands outwards. In the duration of the laser pulse a considerable amount of material is released into the vacuum. The subsequent interaction between the released material and the laser pulse greatly reduces the efficiency of laser energy coupling to the target. However a new class of lasers, whose pulse-widths last hundreds of femtoseconds or less, soon became available which avoided this hydrodynamic expansion problem. In these "ultra-short-" (or simply "short-") pulse laser experiments, the laser energy is deposited before any significant expansion of the target occurs, and results in a more efficient deposition of laser energy to the target. Therefore, using a relatively small short-pulse laser system, it became possible to access regimes of extreme high pressure that are traditionally only attainable using massive long-pulse laser systems. Since the features of short-pulse laser experiments differ considerably from that of long-pulse laser experiments, new approaches for their use in studying strongly coupled plasmas are required.

1.2.1 Long-Pulse Laser Experiments

The process of laser-driven ablation using long-pulse lasers is shown in Figure 1.1. The laser light is focussed onto the target surface, where initially via skin depth deposition the laser energy is absorbed predominately by free electrons. Through electron-ion collisions, the ions are rapidly thermalized, causing the surface layer of the target to be vaporized. As the vaporized material expands away from the surface, a rarified or coronal plasma is formed. As a result a large density gradient, ranging from vacuum to solid density, is formed. Associated with all plasmas is a physical quantity called the plasma frequency,

Figure 1.1: Schematic diagram of a solid target irradiated by a laser at high intensities. The laser is incident from the right. The light gray area is the coronal plasma, the black line is the critical density n_{cr} layer, the dotted area is the ablation zone, and the dark gray area is the shock compressed target material.

which is defined as

$$\omega_p = \left(\frac{4\pi n_e e^2}{m_e} \right)^{1/2} \quad (1.1)$$

where n_e is the electron density, m_e is the electron mass and e is the electronic charge. Within this density gradient lies a particular region with a critical density n_{cr} in which the plasma frequency equals the laser frequency. This region is called the critical density layer. For long-pulse laser experiments, long after the coronal plasma is created, laser energy continues to be absorbed by the plasma, but only up to the critical density layer. Any light which reaches the critical density layer will generally be reflected back and is absorbed by the coronal plasma. Any light which does penetrate beyond the critical density layer will be an evanescent wave. Therefore, in these experiments the laser-matter interactions are dominated by the rarefied plasma, so any optical measurements made on the target would not provide any clear information about the dense plasma of interest.

Beyond the critical density layer lies a transition region between the hot coronal plasma and the cold, solid target material known as the conduction or ablation zone. In the ablation zone a large temperature gradient exists, where thermal conduction transports energy from the hot, coronal plasma to the cold, solid region. As the solid target material is rapidly ablated, it expands outward, and as a result of conservation of momentum, a large ablation pressure is produced near the surface causing a compression wave to propagate into the solid region. As the laser continues to irradiate the target and ablate material, the ablation pressure increases and causes compression waves of increasing magnitude to be launched into the target. The speed of propagation for these wave is given by the local isentropic sound speed c_s :

$$c_s = \left(\frac{\partial P}{\partial \rho} \right)_S^{1/2} \quad (1.2)$$

where P is the pressure, and ρ is the material density. For adiabatic compression the

sound speed can be written as

$$c_s = \left(\frac{\gamma P}{\rho} \right)^{1/2} \quad (1.3)$$

where γ is the ratio of specific heat. It can be shown that for most materials, the sound speed increases with density (for an ideal gas, $c_s \propto \rho^{\frac{\gamma-1}{2}}$ with $\gamma > 1$), therefore the subsequent, stronger waves will catch up with preceding, weaker waves and leads to a steepening of the wave front. Eventually, a sharp discontinuity of the thermodynamic variables between the uncompressed and compressed material is formed. This is the general picture of a shock wave.

1.2.2 Short-Pulse Laser Experiments

In short-pulse laser experiments, the process of driving a shock into a material is similar to that of the long-pulse experiments except for a few key aspects. Because time scales are too short for the creation of a layer of expanded plasma to dominate the laser-matter interaction, more energy is directly transferred to the solid density material. During the extremely brief laser pulse, the energy is absorbed via skin depth deposition almost entirely by the electrons, and due to the finite time required for energy exchange between electrons and ions, the latter initially remain cold. Within a thin solid layer of material at the target surface, the electrons can reach temperatures of over a hundred eV. Consequently, while the electrons begin to thermalize with the ions, a large thermal flux drives a supersonic heat wave into the bulk of the solid target [17, 18]. As the ions are heated, the surface is ablated and material expands into the vacuum. Once again, because of conservation of momentum a compression wave ahead of the ablation front is generated and a shock front is formed. As the thermal wave moves away from the surface it slows and the shock wave soon overtakes it.

In long-pulse laser experiments a relatively steady shock wave is formed, but in short-pulse laser experiments the shock wave is quite transient. At the end of the short laser pulse no more energy is absorbed at the surface. Hence the pressure, density and temperature of the shock wave decrease rapidly resulting in a *decaying* shock wave. Nevertheless, as long as the shock states does not change considerably while measurements are being made, the non-steady nature of the short-pulse laser driven shock wave would not preclude their use in the study of strongly coupled plasmas. However in order to properly make use of the short-pulse laser method, detailed understanding of such decaying shocks and new experimental approaches are required.

1.3 Present Work and Thesis Outline

Because of the transient nature of the shock wave produced by short-pulse lasers, most earlier works [19, 20, 21] have avoided studying the actual shock wave. Instead, they have focussed on the interaction between the femtosecond laser pulse with the near-solid density plasma formed just after the target is irradiated [19, 20]. Nevertheless, issues related to hydrodynamic expansion and plasma scale-lengths still remain pertinent. The use of ultra-thin foils ($\sim$ few hundred Å) along with femtosecond lasers has been suggested as a new approach to produce a well-defined solid density plasma, or *slab plasma* [23].

Recently, optical probing of femtosecond-laser driven shock waves in aluminum were performed by Evans *et al.* [24]. In their work, the shock states were inferred from the simultaneous measurement of the particle and shock velocities along the principle Hugoniot curve, as is normally done in long-pulse laser experiments. However, it will be shown in this thesis that for shock waves driven by femtosecond lasers the shock states cannot be interpreted in the usual manner. The focus of the present work is to delineate the unique characteristics of the shock wave which is produced in short-pulse

laser experiments. Numerical simulations are used to study the features of these decaying shock waves and to analyze a new approach for obtaining information on strongly coupled plasmas, such as the electrical conductivity.

A description of the hydrodynamic model and electromagnetic wave solver used in the calculations is presented in chapter 2. In chapter 3, results from the simulation of a proposed experimental scheme are discussed. A summary of the thesis is provided in chapter 4.

Chapter 2

Modelling of Femtosecond-Laser Driven Shock Waves

Due to the complex nature of laser-matter interactions, computer modelling of laser induced phenomena have become standard. A finite element model in which the plasma is treated as a fluid is used. The *hydrocode* LTC (Laser Target Code) used in the present work was adapted from the laser fusion code MEDUSA [25] in 1987 by Peter Celliers. As various investigations have been pursued, LTC has undergone considerable modifications due to the changing requirements. In this chapter, a brief overview of the operation of the code is presented. A more complete description can be found elsewhere [26, 27].

2.1 Shock Wave Theory

An idealized shock transition between initial uncompressed material and final compressed material is illustrated in figure 2.1. The locus of thermodynamic states obtained from the compression of a material by a single shock wave is uniquely determined by the equation of state (EOS). This locus, usually referred to as the shock adiabat or the Rankine-Hugoniot curve (or simply the Hugoniot), can be calculated in a completely general manner by considering the conservation of mass, momentum and energy across the shock front. The respective mass, momentum and energy conservation equations across the shock front are

$$\rho_0 U_S = \rho_1 (U_S - U_P) \quad (2.1)$$

$$P_1 - P_0 = \rho_0 U_S U_P \quad (2.2)$$

$$E_1 - E_0 = \frac{1}{2}(P_1 + P_0) \left(\frac{1}{\rho_0} - \frac{1}{\rho_1} \right) \quad (2.3)$$

where ρ is the mass density, P the pressure, E the internal energy, and U_S and U_P are the shock speed and the particle speed. The subscripts of 0 and 1 denote the unshocked and shocked regions, respectively. For the shock state, this gives three equations with five parameters: U_S , U_P , P_1 , ρ_1 , and E_1 .

If the EOS, $P = P(\rho, E)$, is known for the material, then it may be added as a fourth equation which reduces the system of equations to a single equation and one free parameter. The resulting single equation gives a locus of points attainable by a single shock for specific initial conditions (ρ_0, P_0, E_0) known as the Hugoniot curve, see figure 2.2. Alternatively, if the EOS of the material is not known it can be measured in shock wave experiments (e.g. U_S and U_P measurements). From the three conservation equations and simultaneous measurement of two shock parameters the EOS parameters P_1 , ρ_1 and E_1 (along the Hugoniot) can be calculated [28].

2.2 Material Models

The EOS along with other material models used in the present work is briefly described. In order to simulate plasma phenomena, the hydrocode requires the equation of state (pressure P and temperature T), average ionization Z^* , thermal conductivity κ and electrical conductivity σ for the material of interest. These are usually given by the

Figure 2.1: Schematic diagram of an idealized shock transition. The particle speed in the uncompressed material is assumed to be negligible.

Figure 2.2: Schematic diagram of Hugoniot curve on the pressure-density plane.

various models as functions of ρ and E ,

$$\begin{aligned}
 T &= T(\rho, E) \\
 P &= P(\rho, E) \\
 Z^* &= Z^*(\rho, E) \\
 \kappa &= \kappa(\rho, E) \\
 \sigma &= \sigma(\rho, E).
 \end{aligned} \tag{2.4}$$

2.2.1 Equation of State and Average Ionization

For this work, a quotidian equation of state (QEOS) based on the work by More *et al.* [29] is used. QEOS is a general purpose equation of state model. It is a self-contained theoretical model that requires no external data base, and the inputs are simply the material composition and the cold solid properties (such as its density, bulk modulus or sound speed). In this model, the electrons and ions are described separately, thus the total internal energy and pressure of the system is, to a sufficient degree of accuracy, assumed to be the sum of the electron and ion contributions. The thermodynamic variables can be expressed in the following form:

$$T_e = T_e(\rho, E_e), \quad P_e = P_e(\rho, E_e) \tag{2.5}$$

$$T_i = T_i(\rho, E_i), \quad P_i = P_i(\rho, E_i) \tag{2.6}$$

The physical properties of the ions are treated completely independent of the electrons except for the constraint that the charge densities are equal. The ion equation of state is based on a Cowan fluid model [30] which combines the Debye, Grüneisen, and liquid scaling-law theories.

The electron equation of state is based on the Thomas-Fermi theory formulated by Feynman, Metropolis, Teller [31] for hot dense plasmas. In the Thomas-Fermi model, the electrons are treated as a charged fluid surrounding the nucleus. The properties of this

electron gas are obtained from finite temperature Fermi-Dirac statistics. Effects due to the plasma environment are introduced using the ion-sphere model where each nucleus is located at the centre of a spherical cavity of radius

$$R_0 = \left(\frac{4\pi n_i}{3} \right)^{-1/3} \quad (2.7)$$

and the cavity contains enough electrons to be electrically neutral while other ions are assumed to remain outside the sphere. The ion density n_i is defined as

$$n_i = \frac{\rho}{Am_p} \quad (2.8)$$

where A is the average atomic mass number, and m_p is the proton mass. In addition, the model is modified by a semi-empirical treatment due to Barnes [32] to account for the chemical bonding energy, so that the pressure at solid density at 300 K is not unrealistically large (several MBar), but only a few kBar. The free electron density n_e is given as

$$n_e = Z^* n_i. \quad (2.9)$$

In QEOS, the Thomas-Fermi theory is also used to obtain the average ionization of the plasma which is assumed to be in local-thermodynamic-equilibrium (LTE) [33, 34]. However, in order for plasma to be considered in LTE, the following conditions must be satisfied: (a) collisional ionization, recombination, excitation, and de-excitation are dominant over all radiative transitions; (b) the electron and ion temperatures are equal; and (c) the electron and ion velocity distributions are near Maxwellian. When these conditions are not satisfied, a non-LTE ionization model such as a collisional-radiative equilibrium (CRE) model [35, 36], or a time-dependent collisional-radiative (CR) model [37, 38], usually has to be employed. The CRE model is used to calculate the average ionization for the case when radiative atomic transitions compete with collisional transitions, but the plasma conditions are changing slowly enough so that the distribution

of ionization states can adapt to the changing plasma conditions. On the other hand, the CR model is used to calculate the case when the plasma conditions are changing faster than the rate at which the ionization state distributions can adapt. However for the present work, the effect of non-equilibrium ionization is not included, thus ionization values used are those given by QEOS.

2.2.2 Thermal and Electrical Conductivities

Thermal and electrical conductivities are described by the conductivity model developed by Lee and More [40]. In this model, the electron transport coefficients are obtained from the solution of the Boltzmann equation in the relaxation time approximation [39]. In this approach, one solves the Boltzmann equation for the electron distribution function and then uses the solution to calculate the flux of electrical current and energy. The electron relaxation time is calculated from the electron collision frequency which only includes contributions from electron-ion and electron-neutral collisions. Calculation of the electron transport coefficients requires the knowledge of the momentum transfer cross sections. In the conductivity model, the Coulomb cross section

$$\sigma_{tr} = \frac{4\pi(Z^*)^2 e^4 \ln \Lambda}{m^2 v^4} \quad (2.10)$$

is employed using appropriate cut-off parameters. The Coulomb logarithm $\ln \Lambda$ is given by

$$\ln \Lambda = \frac{1}{2} \ln(1 + b_{max}^2/b_{min}^2) \quad (2.11)$$

where b_{min} and b_{max} are the upper and lower cut-offs on the impact parameters for Coulomb scattering. The minimum electron-ion impact parameter b_{min} is the value allowed by the uncertainty principle, while the maximum impact parameter b_{max} is the Debye-Hückel screening length [41], corrected for electron degeneracy. In accordance with other numerical results [42] a minimum value of 2 is set for the Coulomb logarithm in

order to overcome the inherent difficulty in the model for the case when the calculated electric field screening length become less than the interionic spacing.

2.3 Fluid Model

2.3.1 Hydrodynamics

LTC solves the one-dimensional Lagrangian fluid equations involving the conservation of mass, momentum, and energy. Initially, the target is partitioned into a mesh of fluid cells. In the Lagrangian description, the calculation follows the time evolution of the cells in which the mass of each cell is kept constant. In this scheme, the size of the cells decrease or increase to respond to the respective processes of compression and rarefaction of the target. The independent variables are time t , and the Lagrangian coordinate of mass m (kg/m^2) which is defined in terms of the density profile $\rho(r, t)$ as

$$m(r, t) = \int_{R_0}^r \rho(r', t) dr' \quad (2.12)$$

where R_0 is a reference position, taken to be zero, and r is the cell position in the laboratory frame. The target material is treated as a compressible fluid composed of electron and ion components each described by a separate temperature. A more complete treatment of the model would include a radiation component, with a corresponding separate temperature, as well. However, in the present work the radiation component is negligible, thus effects due to radiation transport are omitted.

In Lagrangian coordinates the fluid equations are

$$\frac{\partial V}{\partial t} - \frac{\partial u}{\partial m} = 0 \quad (2.13)$$

$$\frac{\partial u}{\partial t} + \frac{\partial (P_e + P_i)}{\partial m} = 0 \quad (2.14)$$

$$\frac{\partial E_e}{\partial t} + P_e \frac{\partial V}{\partial t} = Q_e \quad (2.15)$$

$$\frac{\partial E_i}{\partial t} + P_i \frac{\partial V}{\partial t} = Q_i \quad (2.16)$$

Equation 2.13 represents the fluid continuity equation, where $V = 1/\rho$ is the specific volume, and u is the fluid velocity. Conservation of momentum of the fluid is represented by equation 2.14, where P_e and P_i are the pressures due to the electrons and ions, respectively. Energy conservation in the fluid is represented by the two equations 2.15 and 2.16, where E_e and E_i are the internal electron energy and ion energy, respectively. The heat source terms Q_e and Q_i for the corresponding electron and ion components of the fluid are discussed in the next section.

For each time step in the simulation, the conservation equations are solved in the following manner. First, the mesh coordinates r and velocities u are advanced in time by solving a simplified version of equations 2.13 – 2.16 which is represented by the following set of equations,

$$\frac{\partial V}{\partial t} - \frac{\partial u}{\partial m} = 0 \quad (2.17)$$

$$\frac{\partial u}{\partial t} + \frac{\partial p}{\partial m} = 0 \quad (2.18)$$

$$\frac{\partial e}{\partial t} + \frac{\partial (u\rho)}{\partial m} = 0 \quad (2.19)$$

The total energy of the fluid is defined as

$$e = E_e + E_i + \frac{u^2}{2},$$

and the total hydrodynamic pressure is given as

$$p = P_e + P_i.$$

In addition, the heat source terms, Q_e and Q_i , are eliminated, and the fluid is considered to be of a single temperature. Since gradients exist between adjacent cells, the piecewise

parabolic method [45] is used to interpolate values at cell boundaries from the average values of the cells. At each interface, a Riemann shock tube problem [46, 47] is solved to yield average values for the pressure and velocity which are used to obtain hydrodynamic momentum and energy fluxes. Hydrodynamic motion is described by using the solutions from the Riemann problem at each interface to advance in time the cell coordinates and velocities. However, this yields only an adiabatic solution which must be corrected by solving the full energy equations which include the heat source terms.

2.3.2 Energy Transport

The energy equations 2.15 and 2.16 can be written in the form,

$$C_V \frac{\partial T}{\partial t} + B_T \frac{\partial \rho}{\partial t} + \frac{\partial E}{\partial t} |_{hydro} = Q \quad (2.20)$$

where the coefficients C_V and B_T are defined by the EOS,

$$C_V = \frac{\partial E}{\partial T} |_V, \quad B_T = \frac{\partial E}{\partial \rho} |_T.$$

The third term on the left-hand side of equation 2.20 represents the adiabatic work which was obtained in the above hydrodynamic phase of the calculations, while on the right-hand side, Q represents the heat sources arising from thermal conduction, electron-ion equilibration and laser energy deposition. As in MEDUSA [25], the Crank-Nicholson iterative time-centred differencing scheme [48] and Gauss elimination[48] were used to solve the energy equations. Both the electrons and ions are treated using the same numerical method with the only difference being in their respective heat sources,

$$Q_e = H_e + X_{ei} + A \quad (2.21)$$

$$Q_i = H_i - X_{ei} \quad (2.22)$$

where H_e and H_i are the respective electron and ion thermal fluxes, X_{ei} is the electron-ion energy exchange, and A is the laser energy deposition.

Flux-Limited Thermal Conductivity

Energy transport due to thermal conduction is given as

$$H = \frac{1}{\rho} \nabla \kappa \cdot \nabla T \quad (2.23)$$

where κ is the thermal conductivity and T is the temperature. However, comparison with observations [49, 50, 51] show that this classical Fourier heat flow law 2.23 overestimates the thermal flux in areas of the plasma where the mean free path of the electrons are greater than the thermal gradient scale length of the shock front. Therefore, the non-local thermal flux would cause an unsubstantiated preheat ahead of the shock wave.

To reduce thermal flux where steep temperature gradients exist, a harmonic mean flux-limited [52] thermal conductivity κ is used:

$$\frac{1}{\kappa} = \frac{1}{\kappa_{model}} + \frac{1}{fW} \quad (2.24)$$

where κ_{model} is the thermal conductivity given by the conductivity model, and f is a free parameter called the flux limiter. When applied to equation 2.23, W refers to the heat flux that could arise from free streaming electrons, and is given as

$$W = \frac{\partial x}{\partial T} n_e k_B T \bar{v} \quad (2.25)$$

where x is the position, T is the temperature, n_e is the electron density, k_B is the Boltzmann constant, and $\bar{v}$ is the greater of the mean thermal velocity or the Fermi speed. The theoretical maximum value of f is 0.6 since $0.6n_e k_B T \bar{v}$ corresponds to the free streaming limit which is the maximum heat flux that can be carried by electrons with a Maxwellian velocity distribution.

Electron-Ion Equilibration

Because of the short time scales of high intensity short-pulse laser experiments, thermal equilibrium between the electrons and ions within the plasma cannot be assumed. The

traditional description of thermal equilibration between electrons and ions is

$$\frac{dT_i}{dt} = \frac{(T_e - T_i)}{\tau_{eq}} \quad (2.26)$$

where T_e and T_i are the respective electron and ion temperatures, and τ_{eq} is the equilibration time.

In numerical calculations, it is more convenient to consider the energy transferred between the electron and ion subsystems during a single time step,

$$\Delta E = \frac{c_{Vi}(T_e - T_i)}{\tau_{eq}} \Delta t \quad (2.27)$$

where c_{Vi} is the isochoric ion heat capacity and Δt is the time step. For weakly coupled plasmas $c_{Vi} = 3n_i k/2$ where k is the Boltzmann constant. For weakly coupled and non-degenerate plasmas, τ_{eq} is given by an expression by Spitzer [53],

$$\tau_{eq} = \frac{3}{8\sqrt{2\pi}} \frac{m_i T_e^{3/2}}{(Z^*)^3 n_i e^4 \ln \Lambda} \quad [s] \quad (2.28)$$

where m_i is the ion mass, n_i the ion density, Z^* the average ionization of the ions, e the electron charge, and $\ln \Lambda$ the Coulomb logarithm. Brysk [54] extended the Spitzer formula to derive an expression for degenerate plasmas,

$$\tau_{eq} = \frac{3\pi m_i \hbar^3}{8m_e^2 (Z^*)^3 e^4 \ln \Lambda} (1 + e^{-\mu/kT_e}) \quad [s] \quad (2.29)$$

where m_e is the electron mass, μ the chemical potential, and $\hbar = h/2\pi$ where h is the Planck constant. However, both equations 2.28 and 2.29 for τ_{eq} are valid only for plasmas where the ion-ion correlation is weak. For strongly correlated systems, these expressions underestimate the equilibration times by two to three orders of magnitude compared to observations [55, 56, 57, 58, 59]. Unfortunately, no simple analytical expressions are available for τ_{eq} for strongly coupled plasmas.

The factor c_{Vi}/τ_{eq} is more conveniently replaced by a single constant g which is left as a free parameter. The phenomenological approach of the electron-ion energy exchange is modelled as

$$\Delta E = \frac{\rho}{\rho_0} \frac{g}{\rho_0} (T_e - T_i) \Delta t \quad [J/kg] \quad (2.30)$$

where g is the electron-ion coupling coefficient, and the density dependence of the electron-ion equilibration time is incorporated by scaling the density of the plasma ρ with the solid density ρ_0 of the material of interest.

Laser Energy Deposition

For electrons, in addition to the energy transport due to thermal conductivity and electron-ion equilibration, energy is absorbed from the laser. The mechanics of laser energy deposition is modelled using an electromagnetic wave solver. The plasma is treated as an inhomogeneous dielectric medium with the complex dielectric function,

$$\epsilon(\omega) = 1 + i \frac{4\pi\sigma(\omega)}{\omega} \quad (2.31)$$

where $\sigma(\omega)$ is the electrical conductivity at the laser frequency ω . The Drude model [62] is used to construct the electrical conductivity,

$$\sigma(\omega) = \frac{\omega_p^2}{4\pi(\nu_{ei} - i\omega)} \quad [s^{-1}] \quad (2.32)$$

$$\nu_{ei} = \frac{Z^* n_i e^2}{m_e \sigma_0} \quad (2.33)$$

where ω_p is the plasma frequency, ν_{ei} is the electron-ion collision rate, σ_0 is the tabulated DC electrical conductivity, Z^* is the average ionization, m_e is the electron mass, and $i = \sqrt{-1}$.

The procedure used to solve the wave equations follows treatment given by Born and Wolf [61] for stratified media, and is briefly described in the next section. Laser

energy deposition via Ohmic heating based on the electrical conductivity of the plasma is calculated as

$$A = \langle \mathbf{E} \cdot \mathbf{J} \rangle = \frac{1}{2} \text{Re}(\sigma) |\mathbf{E}|^2 \quad (2.34)$$

where $\mathbf{E}$ is the electric field vector, and $\mathbf{J} = \sigma \mathbf{E}$ is the current density.

2.4 Electromagnetic Wave Solver

In the electromagnetic wave solver, Maxwell's equations are solved using the Helmholtz formulation in the following manner [61]. A medium whose properties are constant throughout each plane perpendicular to a fixed direction is called a *stratified medium*. Consider a plane electromagnetic wave propagating through this stratified medium with the plane of incidence taken to be the yz -plane, as shown in figure 2.3. Any incident electromagnetic wave may be resolved into two linearly polarized orthogonal waves: a transverse electric (TE) wave and a transverse magnetic (TM) wave.

For a TE, or S-polarized, wave incident on the target medium, $E_y = E_z = 0$ and Maxwell's equations reduce to the following scalar equations:

$$\begin{aligned} \frac{\partial H_z}{\partial y} - \frac{\partial H_y}{\partial z} + \frac{i\epsilon\omega}{c} E_x &= 0 \\ \frac{\partial H_x}{\partial z} - \frac{\partial H_z}{\partial x} &= 0 \\ \frac{\partial H_y}{\partial x} - \frac{\partial H_x}{\partial y} &= 0 \\ \frac{i\omega\mu}{c} H_x &= 0 \\ \frac{\partial E_x}{\partial z} - \frac{i\omega\mu}{c} H_y &= 0 \\ \frac{\partial E_x}{\partial y} + \frac{i\omega\mu}{c} H_z &= 0. \end{aligned} \quad (2.35)$$

Figure 2.3: Coordinate system used by the electromagnetic wave solver. The propagation of the plane electromagnetic wave of the laser incident to the stratified medium at an angle θ to the z -axis is shown.

Eliminating H_y and H_z yields the following differential equation

$$\frac{\partial^2 E_x}{\partial y^2} + \frac{\partial^2 E_x}{\partial z^2} + n^2 k_0^2 E_x = \frac{d(\log \mu)}{dz} \frac{\partial E_x}{\partial z} \quad (2.36)$$

where the n is index of refraction of the medium, while k_0 and λ_0 are the respective wave number and wavelength of the incident wave outside the medium, and are given as

$$n^2 = \epsilon\mu, \quad k_0 = \frac{\omega}{c} = \frac{2\pi}{\lambda_0}. \quad (2.37)$$

For all cases in the present work, the magnetic permeability μ is equal to one since only non-magnetic materials are studied. Through separation of variables, equation 2.36 can be written as the following second-order differential equations:

$$\begin{aligned} \frac{d^2 U}{dz^2} - \frac{d(\log \mu)}{dz} U + k_0^2(n^2 - \alpha^2)U &= 0 \\ \frac{d^2 V}{dz^2} - \frac{d[\log(\epsilon - \frac{\alpha^2}{\mu})]}{dz} \frac{dV}{dz} + k_0^2(n^2 - \alpha^2)V &= 0 \end{aligned} \quad (2.38)$$

where

$$\begin{aligned} E_x &= U(z)e^{i(k_0\alpha y - \omega t)} \\ H_y &= V(z)e^{i(k_0\alpha y - \omega t)} \\ H_z &= W(z)e^{i(k_0\alpha y - \omega t)}. \end{aligned} \quad (2.39)$$

The functions U , V and W are related by the following equations:

$$\begin{aligned} U' &= ik_0\mu V \\ V' &= ik_0(\epsilon - \frac{\alpha^2}{\mu})U \\ \alpha U + \mu W &= 0. \end{aligned} \quad (2.40)$$

For the case when the wave is a homogeneous plane wave

$$\alpha = n \sin \theta = \text{constant} \quad (2.41)$$

where θ is the angle which the normal to the wave makes with the z -axis.

Similarly, for a TM, or P-polarized, wave incident on a target medium ($H_y = H_z = 0$), a corresponding set of expressions can be found. As a consequence of the symmetry of Maxwell's equations, the substitution rule of exchanging ϵ and $-\mu$ gives the following equations:

$$\begin{aligned} \frac{d^2 U}{dz^2} - \frac{d(\log \epsilon)}{dz} \frac{dU}{dz} + k_0^2 (n^2 - \alpha^2) U &= 0 \\ \frac{d^2 V}{dz^2} - \frac{d[\log(\mu - \frac{\alpha^2}{\epsilon})]}{dz} \frac{dV}{dz} + k_0^2 (n^2 - \alpha^2) V &= 0 \end{aligned} \quad (2.42)$$

where

$$\begin{aligned} H_x &= U(z) e^{i(k_0 \alpha y - \omega t)} \\ E_y &= -V(z) e^{i(k_0 \alpha y - \omega t)} \\ E_z &= -W(z) e^{i(k_0 \alpha y - \omega t)}. \end{aligned} \quad (2.43)$$

Now, the functions U , V and W are related by the following equations:

$$\begin{aligned} U' &= ik_0 \epsilon V \\ V' &= ik_0 (\mu - \frac{\alpha^2}{\epsilon}) U \\ \alpha U + \epsilon W &= 0. \end{aligned} \quad (2.44)$$

The solutions, subject to the appropriate boundary conditions, of the differential equations (2.38 and 2.42) and various theorems relating stratified medium, can most conveniently be expressed in terms of matrices. In the case of a homogeneous dielectric film, ϵ , μ and $n = \sqrt{\epsilon\mu}$ are constants. For a S-polarized wave, the solutions to 2.38 are given as

$$\begin{aligned} U(z) &= A \cos(k_0 n z \cos \theta) + B \sin(k_0 n z \cos \theta) \\ V(z) &= \frac{1}{i} \sqrt{\frac{\epsilon}{\mu}} \cos \theta \{ B \cos(k_0 n z \cos \theta) - A \sin(k_0 n z \cos \theta) \} \end{aligned} \quad (2.45)$$

with the boundary conditions (at $z = 0$)

$$U(0) = U_0, \quad V(0) = V_0. \quad (2.46)$$

Thus, the x and y components of the electric (or magnetic) vectors in the plane $z = 0$ are related to components in an arbitrary plane z by the *characteristic matrix* $\underline{M}(z)$ of

the medium,

$$\begin{pmatrix} U_0 \\ V_0 \end{pmatrix} = \underline{M}(z) \begin{pmatrix} U_z \\ V_z \end{pmatrix} \quad (2.47)$$

where

$$\underline{M}(z) = \begin{bmatrix} m_{11} & m_{12} \\ m_{21} & m_{22} \end{bmatrix} = \begin{bmatrix} \cos(k_0 n z \cos \theta) & -\frac{i}{p} \sin(k_0 n z \cos \theta) \\ ip \sin(k_0 n z \cos \theta) & \cos(k_0 n z \cos \theta) \end{bmatrix} \quad (2.48)$$

with

$$p = \sqrt{\frac{\epsilon}{\mu}} \cos \theta. \quad (2.49)$$

For a P-polarized wave, the same equations hold, with p replaced by

$$q = \sqrt{\frac{\mu}{\epsilon}} \cos \theta. \quad (2.50)$$

The inhomogeneous target medium can be treated as a stack of thin homogeneous films by multiplying the characteristic matrices of all the stratified media. In other words, the plasma is modelled as a succession of stratified media extending from

$$0 \leq z \leq z_1, z_1 \leq z \leq z_2, \dots, z_{N-1} \leq z \leq z_N$$

with the characteristic matrix

$$\underline{M}(z_N) = \underline{M}_1(z_1) \underline{M}_2(z_2 - z_1) \cdots \underline{M}_N(z_N - z_{N-1}). \quad (2.51)$$

Thus, consider a plane wave incident upon a stratified medium that extends from $z = 0$ to $z = z_l$ and that is bounded on each side by a homogeneous, semi-infinite medium, see figure 2.4

Applying the boundary conditions allow for the calculation of the electric (or magnetic) field strength within the target medium. For a S-polarized wave the boundary conditions are

$$\begin{aligned} U_0 &= E_i + E_r, & U_{z_l} &= E_t, \\ V_0 &= p_0(E_i - E_r), & V_{z_l} &= p_l E_t, \end{aligned} \quad (2.52)$$

Figure 2.4: A schematic diagram of an inhomogeneous medium modelled as a stack of stratified media.

where

$$p_0 = \sqrt{\frac{\epsilon_0}{\mu_0}} \cos \theta_0, \quad p_l = \sqrt{\frac{\epsilon_{z_l}}{\mu_{z_l}}} \cos \theta_{z_l}. \quad (2.53)$$

The amplitudes of the electric field vectors of the incident, reflected and transmitted waves are E_i , E_r and E_t , respectively. Furthermore, ϵ_0 , μ_0 and ϵ_l , μ_l are the dielectric constants and magnetic permeabilities of the first and last media (semi-infinite), while θ_0 and θ_l are the angles which the normals to the incident and transmitted waves make with the z -direction (direction of stratification). The reflection and transmission coefficients of the stratified medium are

$$r = \frac{E_r}{E_i} = \frac{(m_{11} + m_{12}p_l)p_0 - (m_{21} + m_{22}p_l)}{(m_{11} + m_{12}p_l)p_0 + (m_{21} + m_{22}p_l)} \quad (2.54)$$

$$t = \frac{E_t}{E_i} = \frac{2p_0}{(m_{11} + m_{12}p_l)p_0 + (m_{21} + m_{22}p_l)}. \quad (2.55)$$

Thus, the reflectivity and transmissivity, in terms of r and t are

$$R = |r|^2, \quad T = \frac{p_l}{p_0} |t|^2, \quad (2.56)$$

and the phase of the reflected and transmitted waves are

$$\phi_r = \arctan \left[\frac{Re(r)}{Im(r)} \right], \quad \phi_t = \arctan \left[\frac{Re(t)}{Im(t)} \right]. \quad (2.57)$$

For the case of a P-polarized wave the quantites p_0 and p_l are replaced by

$$q_0 = \sqrt{\frac{\mu_0}{\epsilon_0}} \cos \theta_0, \quad q_l = \sqrt{\frac{\epsilon_{z_l}}{\mu_{z_l}}} \cos \theta_l \quad (2.58)$$

in which case r and t are then the ratios of the amplitudes of the magnetic field vectors.

2.5 Other Effects

In high-intensity, short-pulse laser experiments, effects due to radiation (x-rays) transport and suprathermal ("hot") electrons [65] must be addressed. The main concern is the issue

of preheat of the cold material ahead of the shock front by x-rays and/or hot electrons, thus modifying the initial state of the material before it is compressed by the shock wave [66, 67, 68, 69, 70].

Within the thin layer of material heated by the intense short-pulse laser, the electrons may reach extremely high temperatures (100-1000 eV), and emit high-energy photons which penetrate deep into the target. However, for laser irradiances less than 10^{15} W/cm^2 , the electron temperatures of the plasma would be about 100 eV or less, and the power density of the frequency-integrated emission is only about 1 % of the laser radiation even if the plasma radiates as a blackbody. Therefore, the effect of radiation transport would be minimal.

Hot electrons generated in the laser heated layer have large diffusion lengths allowing them to penetrate into the solid and heat the material. Hot electron sources include the processes of spatial dispersion and wave breaking [71] along with the absorption mechanisms of laser energy deposition such as resonant absorption [49], vacuum heating [72], and forward-scattering instabilities [73]. In this thesis, the aforementioned effects are considered negligible since the situations studied involved relatively moderate pump laser intensities ($\Phi_L < 10^{15} \text{ W/cm}^2$) at normal incidence to the target.

Chapter 3

Discussion of Femtosecond-Laser Driven Shock Results

Studying plasma transport properties such as electrical conductivity is of fundamental importance to the understanding of collision processes within a dense plasma. In shock produced plasma experiments, optical probing is a widely used and crucial diagnostic. Studies usually involve the examination of rear surface luminous emission and reflectivity as the shock wave is released at the vacuum-target interface [74]. However, as the release wave expands into the vacuum a critical density layer, which is cooler and of lower density than the shock state, is formed and acts as an absorbing layer to radiation passing through it in either direction. Thus, subpicosecond time resolution is required to allow accurate data collection before the critical density layer expands a few optical depths away from the hot, shock-compressed material.

With the advent of intense femtosecond lasers, a slightly different approach for producing and studying dense plasmas has been developed. The target is irradiated by an intense short-pulse beam to produce near-solid-density plasma with a very short gradient scale-length near the target surface [19]. Because of the limited plasma expansion during the heating pulse, it was thought that the reflectivity of the laser could be calculated by considering an electromagnetic wave propagating through a plasma with a steady-state gradient which is assumed to be linear [20], exponential [20, 21], or given by a power law [22]. From the reflectivity measurements the electrical conductivity of the dense plasma could then be inferred. The initial results of experiments done on aluminum by Milchberg *et al.* [19] using this method showed severe deviations [75] from an existing dense plasma

conductivity model [40]. As a result, several experimental investigations [20, 21] were performed and theoretical models [75, 76] were developed to examine the discrepancy.

However, the apparent discrepancy seemed to hinge heavily on the interpretation of the experimental data. Given the extremely transient nature of the heating pulse, the assumption of a constant gradient scale-length generated by a steady-state ablative flow is questionable. Using a hydrodynamic code to self-consistently treat the plasma dynamics and the laser-plasma interactions, the controversy was resolved by Ng *et al.* [77]. It was shown that the assumption of a single density scale-length throughout the duration of the laser pulse is clearly invalid; it was also revealed that the electromagnetic wave of the laser interacts with the target material over a wide range of plasma conditions. Thus, by incorporating the aforementioned effects, the experimental reflectivity measurements were shown to agree quite closely with the model predictions.

3.1 Proposed Scheme

In order to limit effects due to gradient scale-length issues associated with the aforementioned approach, it is ideally better to look at a shock produced plasma state which is still situated within a material. Traditionally, rear-side optical probing is used to examine the dense plasma state in a shock wave as the shock front approaches the rear surface of a target. Since most studies of dense plasmas deal with materials which are metallic in their cold, solid state, the electromagnetic wave of the probe laser penetrates only a very short skin depth within the material. Hence, observation of the shock-state is quite brief before the shock is released into the vacuum and thereby destroying the desired dense plasma state. To deal with this problem, the rear surface of the target is usually placed on a *window* substrate such as quartz or glass to prevent the shock plasma from releasing, but then effects arising from the interface between the window substrate and the target

material must be addressed. It seems that it would be simpler to study a material with a relatively low absorption coefficient for optical light in the cold, solid state.

At normal conditions, silicon has an indirect band gap of 1.2 eV and its absorption coefficient is $< 10^4 \text{ cm}^{-1}$ for wavelengths $> 0.5 \mu\text{m}$ [78]. Another important characteristic of silicon is that its small band gap closes at a relatively low pressure (~ 12 GPa), thus at higher pressures silicon is expected to be metallic [79] and interband transitions need not be addressed. Silicon also has the practical advantage of being readily available as high-purity, ultrathin wafers with optical-quality surfaces.

A proposed experimental scheme is shown in figure 3.1. On the right, a 400 nm pump pulse with a 120 fs full-width half-maximum (FWHM) pulse-width is incident normally onto the target's front surface. On the left, a 120 fs FWHM, 800 nm probe laser is incident normally onto the target's rear surface. Generally speaking, shorter wavelengths for the pump pulse offer greater absorption and reduce hot-electron production compared to longer wavelengths [69].

A practical target is presented which consists of a $0.1 \mu\text{m}$ thick aluminum layer, and a $1.0 \mu\text{m}$ thick silicon layer situated on a quartz window substrate. In addition, anti-reflective coatings are placed at the window-silicon interface and at the window-vacuum interface. The anti-reflective coatings are needed to prevent interference effects due to probe laser light reflected off the shock front and these surfaces. In the proposed scheme, effects arising from the interface between the window and the silicon are avoided by examining only the results in which the shock wave is still well within the silicon layer. For the present work, the window-silicon interface is considered to be the rear (free) surface of the target.

Because silicon is relatively transparent to optical laser light, the thin aluminum layer is needed to allow for efficient absorption of the pump laser energy in order to initiate and drive a shock wave into the silicon layer. In addition, because the $1/e$ skin depth of

Figure 3.1: Schematic diagram of the proposed scheme. The thicknesses of the anti-reflective coatings and quartz window substrate are not drawn to scale.

aluminum is quite short ($\sim 7nm$ at $\lambda = 400nm$) transmission of the intense femtosecond laser pulse into the silicon target is precluded thereby preventing preheat of the cold silicon. After the aluminum is heated by the pump pulse, a thermal wave propagates into the silicon layer which is then overtaken by a compression wave generated by the ablation of the front surface. As the shock wave moves through the silicon layer, the probe pulse incident on the rear surface propagates through the cold silicon material and interacts with the shock front. Since the probe pulse is extremely short in duration, the measured reflectivity of the shock front is considered almost instantaneous.

For the silicon layer, a thickness of $1.0\mu m$ was chosen for several reasons. First, in contrast to previous studies [80, 81] where thinner targets were used in order to study the thermal wave produced by intense femtosecond lasers, the target used in the present work needed to be thick enough to allow the formation of the shock wave. On the other hand, because there is still a small amount of absorption of the probe laser light by the cold solid silicon, the target cannot be too thick, otherwise the measured reflected probe light will be too attenuated. Therefore, the silicon samples used in this approach would be somewhat restricted to thicknesses $\sim 1\mu m$.

The 120 fs laser pulses employed in the experiment above can be readily obtained from a Titanium:Sapphire ($Ti:Al_2O_3$) laser system [82]. The 800 nm probe pulse would correspond to the fundamental frequency (1ω) beam of the $Ti:Al_2O_3$ laser, while the 400 nm pump pulse, whose irradiance ranges from $10^{14}W/cm^2$ to $10^{15}W/cm^2$, would correspond to the frequency doubled (2ω) beam. In the present work, the temporal profiles of pump and probe pulses are modelled after the laser pulse reported in a recent experiment [83] as shown in figure 3.2. In reality, the wing structure of the pulse extends for a considerable amount of time away from the peak of the pulse, but at such low intensities the target is effected negligibly. For the sake of efficiency, in the simulation the laser pulse is truncated at intensities below 10^{-6} of the peak of the pulse laser intensity.

Figure 3.2: Temporal profile of 120 fs laser pulse used in the simulations.

3.2 Results

3.2.1 Basic Considerations

Recently, work has been done to further shorten the pulse duration down to about 35 fs, and over the last decade many experimental studies have used slightly longer pulse duration (up to ~ 500 fs). For the present work the choice of pulse duration is not too critical. As long as the pulse is short enough to avoid hydrodynamic expansion during the heating of the target by the pump pulse, efficient laser energy deposition would be achieved. In addition, the probe pulse duration should be short enough such that the shock state of the changes negligibly during the interaction with the probe pulse, so it may be considered to be an instantaneous measurement of a shock state. Slight changes in the pump pulse duration have the effect of simply decreasing or increasing the amount of energy that the target absorbs, but adjustment of the pump pulse duration represents a rather inefficient means of laser energy control. A more suitable control over laser energy deposition would be the varying of the irradiance of the pump pulse while keeping the pulse duration constant.

The pump irradiances used in the present work were chosen for several reasons. They were limited to a maximum value of 10^{15}W/cm^2 in order to avoid preheat of the uncompressed material in front of the shock by radiation transport and/or hot electrons. Meanwhile, a minimum value of 10^{14}W/cm^2 was employed in order to avoid phase transition issues. In addition to the fact that silicon become metallic at 12 GPa [79], the sequence of silicon phase transitions include cubic diamond to body-centre tetragonal at 11 GPa [84, 85, 86], to simple hexagonal at 13-16 GPa [87], to an intermediate phase of undetermined structure at 34 GPa, to hexagonal-close packed at above 40 GPa [88], and to face-centre cubic above 78 GPa [89]. Plasmas produced by lower laser irradiances

would have shock pressures comparable to that associated with the these phase transitions, thus complicating matters. For laser irradiances above 10^{14}W/cm^2 , the shock pressures are substantially above these phase transition pressures. Also, for irradiances less than 10^{14}W/cm^2 , the initial shock wave is relatively weak ($\sim$ a few hundred kbars) and does not propagate deep into the material before the thermal wave actually catches back up to the shock wave and extinguishes the shock structure.

Although a harmonic-mean flux limiter model [52] is incorporated in the calculations, the heat flux in the absorption region or the shock front never reaches saturation even for a flux limiter of 0.03 because of the large collision frequency and the limited temperature difference across the heat or shock fronts. In all calculations, a flux limiter of 0.6 was used, thereby avoiding legislation of heat flow.

Since the electrons and ions are treated separately in the calculations and there exists no simple analytical expression for electron-ion energy exchange, the thermal equilibration between the electrons and ions was treated using a phenomenological coupling constant. Based on previous studies by Ng *et al.* on silicon [43], an electron-ion coupling coefficient of $g = 10^{16} \text{W/m}^3 \text{K}$ was used for all calculations in this thesis. Changes in g up to an order of magnitude showed only slight changes in the electron and ion temperatures at the shock front.

3.2.2 Characteristics of the Decaying Shock Wave

Off-Hugoniot States

In shock physics, thermodynamic information about the shock state is usually described by the Rankine-Hugoniot curve (or Hugoniot). For given initial conditions, the final state reached due to compression of the initial state by a single shock will lie on the Hugoniot. For a spatially and temporally constant shock propagating through a sample

all compressed material behind the shock front will be in the same shocked state. In order for a constant shock wave to continue to propagate, a steady or quasi-steady energy source must drive the shock, as in the case of long-pulse laser experiments. Obviously, in short-pulse laser experiments the shock wave is not driven continuously, so it decays in magnitude as it propagates through the sample. Because of the influence of the initial thermal wave in the short-pulse laser experiments, the shock states of the decaying shock wave will lie off the Hugoniot. Figures 3.3 and 3.4 show the pressure and electron temperature at the shock front of a decaying shock wave driven by an intense femtosecond laser in silicon, along with the corresponding silicon Hugoniot states.

The end of the pump pulse is represented by the point **A** where 99.99% of the laser energy has been deposited in the target. At this point in time, the ions at the front surface of the target are still relatively cool while the electrons are heated to their maximum temperature thus producing a very large thermal pressure at the surface of the target. From point **A** to point **B**, a thermal wave propagates into the silicon followed by a compression wave, where at point **B** the compression wave has overtaken the thermal wave and steepened to form a shock front. After point **B**, the magnitude of the shock rapidly decays as the front surface releases into vacuum. Initially, the state at the shock front is well off the Hugoniot due to the effect of the thermal wave on the material. As the decaying shock propagates further into the target, the state at the shock front can be seen to asymptotically approach the Hugoniot states. To obtain more information about the femtosecond-laser driven shock, careful examinations of the thermodynamic spatial profiles of the decaying shock at different times are needed.

Spatial Profiles of the Decaying Shock Wave

Figures 3.5 - 3.12 represent the spatial *snapshots* of the thermodynamic variables of pressure, density, electron temperature and ion temperature at different times. Examination

Figure 3.3: The Si off-Hugoniot states on the pressure-compression plane for a femtosecond-laser driven shock wave. The solid density of Si is $\rho_0 = 2.33 \text{ g/cm}^3$. Time proceeds from point A to B to C.

Figure 3.4: The Si off-Hugoniot states on the electron temperature-compression plane for a femtosecond-laser driven shock wave. The solid density of Si is $\rho_0 = 2.33 \text{ g/cm}^3$. Time proceeds from point A to B to C.

of these profiles will illustrate the unique evolution and characteristics of shock waves generated by intense femtosecond lasers. Results for a pump irradiance of $5 \times 10^{14} \text{ W/cm}^2$ is used as the standard case for discussion. For other pump irradiances, the shape of the profiles are quite similar to that of the standard case while the magnitudes of the profiles scale with the irradiance used. The features of the decaying shock wave are revealed by studying the profiles at early times shortly after the end of the pump pulse, figures 3.5 - 3.8, and at later times when a definite shock wave has been formed, figures 3.9 - 3.12. Time zero corresponds to the time of the peak of the pump pulse laser intensity that is incident on the target.

Referring to the early time profiles, figures 3.5 - 3.8, the highest pressures occur at the front of the target while it is being heated by the pump pulse. At this time, no compression of the target has taken place, and this high pressure is attributed to the large thermal pressure of the electrons heated by the pump laser. Since only the electrons absorb energy from the laser, the electron temperature at the surface reaches up to about 70 eV while the ions remain relatively cold. As the electrons start to thermalize the ions, the surface of the target begins to be ablated causing material to expand into the vacuum. Meanwhile, an electron thermal wave [81, 90] begins to propagate through the target. The pressure at the target front surface quickly decays due to both the release of the hot material into the vacuum and the propagation of the thermal wave into the target. Because no additional energy is provided for the thermal wave as it moves further into the target, it quickly begins to weaken and slow down. At the same time, because of conservation of momentum, the rapidly expanding ablated material leads to compression of the material near the surface.

The compression wave follows the thermal wave into the target and as it moves from the aluminum layer to the silicon layer, its density profile starts to steepen into a discontinuity and a shock wave is formed. The shock wave overtakes the initial thermal

Figure 3.5: Early time pressure spatial profiles. The position $x = 0$ represents the initial location of the target front surface, and the pump pulse is incident from the right. For early times, the profiles remain near the front surface and the material not shown ($x < -0.5\mu\text{m}$) is still cold, solid Si.

Figure 3.6: Early time electron temperature spatial profiles. The position $x = 0$ represents the initial location of the target front surface, and the pump pulse is incident from the right. For early times, the profiles remain near the front surface and the rest of target material which is not shown ($x < -0.5\mu\text{m}$) is still cold, solid Si.

Figure 3.7: Early time ion temperature spatial profiles. The position $x = 0$ represents the initial location of the target front surface, and the pump pulse is incident from the right. For early times, the profiles remain near the front surface and the rest of target material which is not shown ($x < -0.5\mu\text{m}$) is still cold, solid Si.

Figure 3.8: Early time density spatial profiles. The position $x = 0$ represents the initial location of the target front surface, and the pump pulse is incident from the right. For early times, the profiles remain near the front surface and the rest of target material which is not shown ($x < -0.5 \mu\text{m}$) is still cold, solid Si.

wave within several picoseconds, but due to the finite equilibration time between the electrons and ions an definite two-temperature shock structure is produced. As the electrons are thermalizing the cold ions, electron thermal flux across the shock front heats the electrons ahead of the shock wave. As a result an inherent electron thermal *foot* [44] of heated electrons in front of the shock wave is produced. The less coupled the electrons and ions are to each other, the greater the electron thermal flux thus the larger the electron thermal foot. Due to the release of the front surface, a rapid tailing off of the pressure and density of the material behind the shock front is seen. Meanwhile, the decrease in the temperature of the ions behind the shock front is retarded by the electrons as thermal equilibration between the two species continues.

Referring to the late time profiles, figures 3.9 - 3.12, the shock wave can be seen propagating through the silicon layer of the target, and just like the thermal wave, it is seen to weaken and slow down. In addition to the decrease in the magnitude of the shock wave, a relaxation in the gradients at the shock front can be observed. The electron temperature profile shows the most noticeable smoothing of the gradient while the other profiles show less pronounced relaxation effects. As the electron temperature profile relaxes, its gradient scale length increases which causes the electron temperature foot ahead of the shock to become larger; in other words more electrons ahead of the shock are heated. Consequently, ahead of the shock more energy is transferred from the electrons to the ions. As a result, a noticeable ion temperature foot begins to appear ahead of the shock front, too. At earlier times, when a strong shock wave with a steep shock front is observed, almost all of the material in front of the shock wave is described as cold silicon. However the appearance of the ion temperature foot modifies that view, since material ahead of the shock front which has ion temperatures above the melting temperature of silicon would have to be considered as liquid silicon. Therefore as time goes on, the notion of the idealized discontinuity between a cold solid and a hot, dense

plasma becomes less realistic.

Also, whereas in early times the material behind the shock front is seen to drop off very steeply, for later times it appears that the *tail* behind the shock front has relaxed as well. Finally, at even later times (not shown in the figures) the shock wave reaches the rear surface of the target. However, in order to neglect effects due to interaction of the shock wave with the rear edge of the target, only times when the shock wave is still relatively far away from the rear surface are pertinent to this work.

Accordingly, based on the observed features of the spatial profiles of the thermodynamic variables at various times, the following remarks can be made. First, some time is required before a definite shock wave is formed. Hence, probing of the target at too early a time would reveal information mainly about the electron thermal wave structure instead of about the desired shocked states. On the other hand, probing of the target too late a time would be on a relaxed shock wave where results are effected by the melted material ahead of the shock front. Consequently, in order to avoid the early and late time effects, study of the shock states produced by an intense femtosecond laser within a certain time frame in between would be optimal. This is not too much of a concern since the probe pulse is also of femtosecond time duration, hence a fair amount of data could be attained during this optimal time window.

Temporal History of Thermodynamic State at the Shock Front

As previously mentioned, for an idealized shock wave, the shock front represents the discontinuity between the thermodynamic states of the material ahead of and behind the shock wave. All material ahead of the shock front is described by a single thermodynamic unshocked state, while all material behind the shock front is described by a single thermodynamic shocked state. In reality, no matter what method is used to drive a shock wave, gradients in the thermodynamic variables, with associated scale-lengths, always

Figure 3.9: Late time pressure spatial profiles. The position $x = -1.1 \mu\text{m}$ represents the location of the rear free surface of the target (window-silicon interface), and the probe pulse is incident from the left.

Figure 3.10: Late time electron temperature spatial profiles. The position $x = -1.1\mu\text{m}$ represents the location of the rear free surface of the target (window-silicon interface), and the probe pulse is incident from the left.

Figure 3.11: Late time ion temperature spatial profiles. The position $x = -1.1\mu\text{m}$ represents the location of the rear free surface of the target (window-silicon interface), and the probe pulse is incident from the left.

Figure 3.12: Late time density spatial profiles. The position $x = -1.1\mu\text{m}$ represents the location of the rear free surface of the target (window-silicon interface), and the probe pulse is incident from the left.

exists. In addition, the material behind the shock front is never in a single, shock state since the compressed material asymptotically settles towards equilibrium. To further complicate matters, in the case of femtosecond-laser driven shock experiments, because the shock wave is not being continuously driven, it is followed by the release wave which results in the rapid relaxation of the shocked state behind the shock front. Thus, as soon as the target material is compressed to a strongly coupled plasma state of interest, it quickly relaxes into an undesirable hot expanded plasma state.

However, one of the key features about the proposed scheme relies upon interaction of the probe pulse with the shock front. In fact, as it will be shown in the next section, the probe laser light propagates a very short distance beyond the shock front, thus only the desired dense plasma states near the shock front are actually measured.

Referring to the previous figures (3.5 - 3.12), at each instant in time the shock wave is described by the various thermodynamic profiles. A shock wave is generally interpreted as a supersonic compression wave, thus identification of the shock front of the decaying shock wave relies mainly on the examination of the density profiles. To simplify matters, in this work the shock front is represented by the point within each density snapshot at which maximum compression of the material occurs. The corresponding pressure, electron temperature, ion temperature, and density represent the state of the shock front at that instant in time. Figures 3.13 - 3.16 track these thermodynamic variables in time as the shock wave propagates through the target material.

Because of the aforementioned criterion of maximum compression utilized for identifying the shock front, for early times the values given correspond to the thermal wave and the steepening of the compression wave. The shock wave is considered to be fully formed near the time at which the density of the "shock front" reaches the maximum compression. Once the shock wave is formed the common feature about the temporal histories is that all the thermodynamic variables decay in time. In addition, there is a

Figure 3.13: Decay of the pressure at the shock front.

Figure 3.14: Decay of the electron temperature at the shock front.

Figure 3.15: Decay of the ion temperature at the shock front.

Figure 3.16: Decay of the density at the shock front.

striking similarity between the decay of the shock front pressure and electron temperature. This reveals that for femtosecond-laser driven shocks, the pressure of the shock front has a strong electron temperature dependence. Meanwhile, the aforementioned process of continual thermal equilibration behind the electrons and ions is indicated by the slightly less rampant drop off in ion temperature at the shock front.

However, whereas the pressure, electron and ion temperatures experience rapid rates of decline, the density of the shock front decreases rather slowly. Within a certain window in time, from 10 ps to 30 ps after the pump pulse, a comparison between the relative change in the electron temperature and the density provide useful information about the decaying shock wave. Whereas the electron temperature drops by more than 50% from its initial value, the density decreases only by about 10%. Consequently, proper measurements of the decaying shock wave could be used to scan the temperature dependence of dense plasma properties along an approximately constant density region. More specifically, the following section demonstrates how optical probing of the decaying shock wave could be used to assess the electrical conductivity of a strongly, coupled plasma.

3.2.3 Optical Probing

In chapter 2, laser energy deposition to the target material by the pump pulse is modelled using an electromagnetic (EM) wave solver. Similarly, the interaction of the probe pulse with the decaying shock wave is described using a post-processor version of the EM wave solver. In the latter case, because the energy of the probe pulse is considerably less than the pump pulse, the target material is not modified by the probe pulse. Therefore, the hydrodynamic profiles generated by the hydrocode for each instant in time can be directly converted into the corresponding complex dielectric function profiles and used to calculate the reflectivity, transmission and phase change of the probe pulse.

As previously mention in chapter 2, the dielectric function is calculated via the Drude prescription using electrical conductivities based on the model given by Lee and More [40] which is valid for a wide range of plasma conditions. However for the unshocked material ahead of the shock wave, and for the material at the foot of the shock front, the model gives values for the dielectric function which differ considerably from that of reported values. This is to be expected since in the former case the material is still just cold, solid silicon, while in the latter case the material is the uncompressed silicon which has been melted by the thermal foot just ahead of the shock front. In both cases, the dense plasma conductivity would not be valid and the associated dielectric function would be incorrect, hence the dielectric values for solid silicon [91] and liquid silicon [92, 93] (See Table 3.1) are substituted in their place. It should be noted that no relevant temperature dependent dielectric values for solid and liquid silicon are given in the literature. Therefore, regardless of the temperature of the solid silicon, the dielectric value at room temperature (298 K) is used, and regardless of the temperature of the liquid silicon, the dielectric value at the melting temperature (1687 K) is used.

Solid Si ($T_i = 298K$)	Liquid Si ($T_i = 1687K$)
$\epsilon = (13.601, 0.0443)$	$\epsilon = (-21.48, 45.01)$

Table 3.1: Dielectric values for solid and liquid silicon for 800 nm light.

The interaction of the probe pulse with the shock wave is best presented by overlaying the electric field strength of the probe laser on the corresponding dielectric functions of the shock wave, see figures 3.17 and 3.18. The probe laser is incident from the left while the shock wave moves within the target from right to left. Within the cold, solid silicon, the electromagnetic wave of the probe light propagates with little attenuation. However, once probe light reaches the shock front, penetration of the electromagnetic wave is very

Figure 3.17: Penetration of the electric field strength of the probe laser beyond the shock front. The normalized E^2 for the 800 nm probe light at 20 ps after the peak of the pump pulse is shown.

Figure 3.18: Dielectric functions at the shock front. The real part of ϵ (solid line), and the imaginary part of ϵ (dashed line) of the silicon target for 800 nm light at 20 ps after the peak of the pump pulse is shown.

short, hence any interaction with the shocked material behind the shock front would be rather weak. If this is the case, then the interface between the cold, solid silicon and the shock front of the decaying shock wave may be approximated as a classical Fresnel surface.

Fresnel Reflectivity Approximation

In most experimental endeavors, the measurement of a definite or single state is usually desired. Since the probe light does not “see” very far beyond the shock front, an approximation may be made where the effect of the shock states in the tail of the decaying shock wave is neglected. To compare the discrepancy between the results given by the EM wave solver on the decaying shock wave and the results for the case of the idealized single shock state the following approximation is used.

For each instant in time the shock profile is simply divided into two regions: (1) a non-conducting, linear medium and (2) a conductor. The boundary between these two regions is given by the shock front of the shock wave. Recall that the location of the shock front is defined as the point in which maximum compression of the material occurs. The density and electron temperature of that point are used to obtain the corresponding dense plasma electrical conductivity which, via the Drude model, gives a dielectric value. Since the probe laser does not penetrate much beyond the shock front, all material behind is assumed to have the same dielectric value. Meanwhile, all material ahead of the shock front is assumed to have the dielectric value corresponding to that of cold, solid silicon, thereby neglecting the effects due to gradient of the shock front and the melted material ahead it.

Under these approximations, for each instant in time the interaction of the probe light with the shock wave is reduced to a Fresnel surface problem that is easily solved using the Fresnel equations. For normal incident light, the Fresnel reflectivity and transmittivity

[94] are given as

$$R_F = \left| \frac{1 - \beta}{1 + \beta} \right|^2, \quad T_F = \left| \frac{2}{1 + \beta} \right|^2 \quad (3.1)$$

where

$$\beta = \frac{\mu_1 n_2^*}{\mu_2 n_1^*} = \frac{n_2^*}{n_1^*} = \frac{\sqrt{\epsilon_2}}{\sqrt{\epsilon_1}}, \quad \mu_1 = \mu_2 = 1. \quad (3.2)$$

The respective media complex refractive indices n_1^* , n_2^* and complex dielectric values ϵ_1 , ϵ_2 are related by

$$n^* = n(1 + i\kappa) = \sqrt{\epsilon} \quad (3.3)$$

where n is the real refractive index and κ is the attenuation index or extinction coefficient.

Recall that a wavelength of 800 nm was chosen for the probe laser because the absorption at that wavelength in cold, solid silicon is minimal. Nevertheless, some attenuation of the probe light will occur as the incident electromagnetic wave travels through the cold material to the shock front, and then as the reflected electromagnetic wave travels back through the cold material to the free surface. For a plane, time harmonic wave, the energy density w of the wave is proportional to the time average of $\mathbf{E}^2$, so it follows that w decreases in accordance with the relation [61]

$$w = w_0 e^{-\chi(\mathbf{r} \cdot \mathbf{s})} \quad (3.4)$$

where χ is the absorption coefficient given by

$$\chi = \frac{2\omega}{c} n\kappa = \frac{4\pi}{\lambda_0} n\kappa \quad (3.5)$$

and λ_0 is the probe laser wavelength in the vacuum. Applying the above attenuation effect to the calculated Fresnel reflectivity of the shock front, the anticipated reflectivity observed at the free surface would be

$$R = R_F e^{-2\Delta x \chi} \quad (3.6)$$

where Δx is the distance from the free surface to the shock front at any given instant in time as seen in figure 3.19. It should be noted that the calculation of transmittivity using the Fresnel approximation would be invalid since behind the shock front the expanded plasma would absorb almost all light propagating beyond it, so only reflectivity results are analyzed.

For early times when the shock wave is not fully formed yet, the Fresnel approximation would not be valid, hence those results are not included. After the shock wave is formed, it can be considered a moving mirror whose refractive index changes with time. As time goes on, the reflectivity of the moving mirror steadily seen at the free surface increases due to two effects. First, as time passes the shock wave moves further into the silicon target so there is less material between the shock front and the free surface which results in less attenuation of the rear probe light. However the increase in the reflectivity is mainly due to the change in the plasma state at the shock front as the shock wave decays. For the dense plasma model used, the electric conductivity at the shock front steadily increases as both the density and electron temperature at the shock front decreases, thus the observed increase in reflectivity.

The above approximation is used for comparison with the result given by the EM wave solver on the hydrodynamic profiles, see figure 3.20. In the latter case, for each plasma state within the decaying shock wave the corresponding dielectric value is used and the effect due to melted silicon ahead of the shock front is included. The following remarks can be stated about the reflectivity of the decaying shock wave calculated by the EM wave solver.

First, at the peak of the pump pulse the boundary between the aluminum pusher layer and the silicon target layer is still intact, thus we see the corresponding reflectivity due to an aluminum-silicon interface. As the thermal wave and then the compression wave moves from the aluminum into the silicon, the reflectivity drops until the shock wave

Figure 3.19: Fresnel reflectivity at the shock front and at the free surface.

Figure 3.20: Reflectivity at $\Phi = 5 \times 10^{14} \text{ W/cm}^2$: Fresnel and EM wave solver. The dotted line is the reflectivity calculated using the Fresnel approximation. The solid line is the reflectivity calculated using the EM wave solver where effects due to the thermal foot and interaction of the probe light beyond the shock front are included.

begins to form. Once the shock wave begins to form and propagates into the silicon, the reflectivity is seen to increase with time as anticipated. It is observed that while the reflectivity given by the EM wave solver is always lower than the reflectivity calculated using the Fresnel approximation, the amount of discrepancy varies with time. Initially, the wave solver reflectivity is lower than the Fresnel reflectivity by a fair amount, as time goes on the wave solver reflectivity begins to approach the Fresnel reflectivity. However, at later times the wave solver reflectivity diverges from the Fresnel reflectivity as expected based on analysis of the spatial profiles.

The above behaviour of the reflectivity may be better understood by directly examining the probe laser interaction with the decaying shock wave. In other words, the electric field strength of probe laser and the complete dielectric functions of the shock structure is examined. Using the wave solver, the probe light penetration of the probe light beyond the shock is seen to remain fairly constant ($\sim 80 \text{ \AA}$) as a function of time, see figure 3.21. Therefore, the deviation of the wave solver reflectivity from the Fresnel reflectivity is traced to changes in the dielectric functions of the melted silicon foot ahead of the shock front and the silicon plasma material behind it as time progresses, see figures 3.22 - 3.24.

At early times, right after the shock front is formed, the dielectric function of the thermal foot is rather small in length so its effect on the reflectivity would be minor. However, behind the shock front the dielectric values of the plasma material differ significantly from that of the shock front. As the shock wave moves further into the target, the dielectric function behind the shock front is seen to level off somewhat, resulting in the wave solver reflectivity to move closer to the Fresnel reflectivity. As time goes on, the thermal foot grows larger and its effect becomes prominent. Also, the relaxation of the shock front is seen in the late time dielectric functions. These two effects are responsible for the divergence of the Fresnel reflectivity and the EM wave solver reflectivity at late times.

Figure 3.21: Penetration depth of probe light beyond shock front. The penetration depth $d_{1/e}$ is defined as the distance beyond the shock front in which the normalized E^2 of the probe light has fallen to $1/e$ of the value at the shock front. Note: at 5 ps the shock front has not steepened entirely thus the larger penetration depth.

Figure 3.22: Dielectric functions at shock front (at 10 ps). The solid line is the real part of ϵ and the dashed line is the imaginary part of ϵ

Figure 3.23: Dielectric functions at shock front (at 20 ps). The solid line is the real part of ϵ and the dashed line is the imaginary part of ϵ

Figure 3.24: Dielectric functions at shock front (at 40 ps). The solid line is the real part of ϵ and the dashed line is the imaginary part of ϵ

Phase Shift

While the above discussion dealt only with the measurement of the magnitude of the reflected probe laser light, the following deals with the phase shift Φ of the reflected probe laser light. The phase shift is defined as

$$\Phi = \Phi_i - \Phi_r \quad (3.7)$$

where Φ_i and Φ_r are the incident and reflected phase of the probe light. The change in the phase of the probe light can be attributed to the following effects: (1) changes in the optical properties of the reflecting surface, (2) changes in the refractive index of the material between the reflecting surface and the observer, and (3) the motion of the reflective surface [95].

In order to analyze the phase shift results, the following approximation is made. The decaying shock wave is treated once again as a moving mirror where the associated motion of the shock front results in Doppler phase shifting which dominates the other two effects. If this is the case, then the phase shift of the reflected probe light would provide a measurement of the shock speed U_S of the decaying shock wave. For the simple case of a mirror moving, its speed is calculated as

$$U_S = \frac{\lambda_0}{4\pi n} \frac{\Delta\Phi}{\Delta t} \quad (3.8)$$

where n is the refractive index of the medium (cold silicon) which the mirror is moving through, λ_0 is the probe light wavelength in the vacuum (800 nm), and $\Delta\Phi$ is the change in the phase shift in a given Δt time interval. The time history of the phase shift of the probe laser light is presented in figure 3.25. Initially, the phase shift remains constant as the probe laser reflects off only the aluminum-silicon interface. The phase is observed to change drastically once the thermal wave reaches the silicon layer and continues to increase as the shock wave propagates through the silicon. The phase shift is seen to

Figure 3.25: Phase shift.

Figure 3.26: Shock speed of decaying shock wave. The solid line is the calculated shock speed using the hydrodynamic profiles, and the dashed line is the calculated shock speed using the phase shift of probe laser.

slowly curve over indicating the decrease in shock speed as expected for the case of a decaying shock wave.

The shock speed is calculated using equation 3.8 and compared with the shock speed calculated directly using the hydrodynamic profiles, see figure 3.26. It is seen that the phase shift calculated shock speeds are in relatively close agreement with that of actual shock speeds. In shock wave experiments, measurement of the shock speed along with measurement of the particle speed are used in conjunction to obtain an EOS measurement. However, as stated earlier the decaying shock wave initially lies well off the Hugoniot and asymptotically approaches it. Hence for femtosecond-laser driven shocks, measuring the shock speed would not allow any direct quantitative access to the actual state of the shock front, rather it may be used as a qualitative check on the steepness of the shock front, thus allowing the use of the moving mirror or Fresnel surface approximation. The state of the shock front would still have to be attained through computer simulations based on the experiment parameters and plasma models.

Electrical Conductivity

The proposed scheme is supposed to be applicable within the pump pulse laser irradiances of 10^{14}W/cm^2 to 10^{15}W/cm^2 . The wave solver results for the two irradiance limits are shown along with the standard case, see figure 3.27. The different irradiances produce shock waves of corresponding magnitudes and shock speeds as seen in the different times in initial drop in reflectivities signifying the their emergence into the silicon. However, once the shock waves are formed, not only do the observed reflectivities follow the same rate of increase, but the magnitudes of the reflectivities have very similar values.

This does not necessarily mean that the proposed scheme lacks the capability to distinguish between the obvious different plasma states which would be produced by the various irradiances. Rather it has to do with the fact that in this regime the dense

Figure 3.27: Similar reflectivity at different irradiances.

Figure 3.28: Similar σ_0 given by the Lee and More model at different irradiances.

plasma model used give electrical conductivities which are very similar, see figure 3.28, thus producing the similar reflectivities seen.

Given that only one electrical conductivity model for silicon was available, further discussion about the optical probing results is somewhat limited. However, as a final remark, the sensitivity of the electrical conductivity on the magnitude of the reflected probe light is discussed. To simplify matters, the Fresnel reflectivity is used to examine the change in reflectivity based on the change in electrical conductivity at the shock front only. Suppose that, at the shock front, the plasma state remains the same but the actual associated electrical conductivity is a factor of 2 higher or lower than that predicted by the model. The effect on the reflectivity would be rather significant ($\sim 10\%$) and would be quite measurable, as seen in figure 3.29. If another electrical conductivity model is provided then the proposed scheme could be used to provide a method for distinguishing between models.

Figure 3.29: Sensitivity of electrical conductivity at the shock front on the Fresnel reflectivity at the free surface. ($\Phi = 5 \times 10^{14} \text{ W/cm}^2$) The solid line represents the case when electrical conductivity σ_0 given by the Lee and More model is used, the dotted line represents the case when $2 \times \sigma_0$ is used, the dashed line represents the case when $\sigma_0 \div 2$ is used.

Chapter 4

Conclusions

4.1 Summary

Using numerical simulations, the characteristics of a shock wave generated by an intense femtosecond laser were investigated. A new approach was developed to study the electrical conductivity of silicon in the dense plasma state. The unique experimental scheme involved the use of a 400 nm, 120 fs pump laser pulse as the shock driver and a 800 nm, 120 fs probe laser pulse as the probe.

The new understanding of the crucial characteristics of femtosecond-laser driven shocks points to an existing controversy in the literature [24] in which the effects of the thermal wave were not addressed and the resulting off-Hugoniot states were not recognized. It was shown that the shock states of shock wave generated by an intense femto-second laser cannot be described by the use of the principle Hugoniot curve as in previous long-pulse laser experiments.

Spatial profiles of the thermodynamic variables at various times showed the resulting shock wave to be quite transient. In addition, the temporal histories of states at the shock front revealed that although the pressure, electron temperature and ion temperature decayed quite rapidly, the mass density at the shock was shown to decrease at a considerably slower rate. With this information, we suggest that the transient nature of the femtosecond-laser driven shock can be used advantageously. Namely, a temperature scan of the electrical conductivity of dense plasma at a relatively steady density can be

performed by optically probing the decaying shock wave as it propagated inside silicon.

The concept of measuring definite shock states as a function of time was addressed using a Fresnel surface approximation. Within a relatively long time window the approximation appears to be reasonably valid and state of the shock front may be discriminated from the other shock states of the decaying shock wave.

The results of this thesis has recently been presented at the International Workshop on Warm Dense Matter [96] and the IEEE International Conference on Plasma Science [97].

4.2 Future Work

Another electrical conductivity model based on density functional theory (DFT) [98, 99] has been constructed by Perrot and Dharma-wardana [76]. In this model all the ingredients needed for the calculation are computed from first principle in a self-consistent manner, using a density functional description of the electrons and ions in the plasma. For aluminum [27], in some regimes the electrical conductivities given by the DFT model have been shown to differ noticeably from the results given by the Lee and More model. Acquisition of electrical conductivities calculated using the DFT model for silicon is being pursued. Using the new electrical conductivities and employing the proposed scheme, reflectivity calculations will be performed. These results may then be compared with the results given in this thesis, thereby allowing for a comparison between the two conductivity models.

For this work, optical probing of the decaying shock wave was limited to a normally incident probe pulse on the rear surface of the target. As a result, only forecasting calculations were possible, that is, the predicted reflectivity of the probe is compared with the measured value. This allows the inference of electrical conductivity through

the numerical simulation. In future work, one may use various angles of incidence of the probe. If reflectivity measurements for both S-polarized and P-polarized light are performed then backcasting calculations would be possible. Using the simultaneous reflectivity measurement of both S-polarized and P-polarized light, the dielectric functions of the shock state may be directly calculated to yield the electrical conductivity of the plasma state. This derived electrical conductivity can then be compared to that given by conductivity theories.

Other materials may also be studied by using slightly modified version of this approach. For instance, carbon, in its various forms, is a material of great interest. It, too, becomes metallic above certain pressures, but the actual features of the transition from insulator to conductor are not well known. By using the decaying shock wave and optical probing, the pressure dependence insulator-conductor transition may be mapped out.

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