SHOCK WAVE AS A PROBE OF FLUX-LIMITED THERMAL TRANSPORT IN LASER-HEATED SOLIDS

By

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Abstract

Classical Spitzer-Härm treatment of electron thermal conduction breaks down in the steep temperature gradients found in laser-heated solids. A phenomenological flux-limiter, which reduces heat flow, is incorporated into computer codes which model laser-target interactions. There is disagreement in what the correct value of the flux-limiter, $f$, should be. A simple method to determine the best value is presented. It involves comparing experimental shock speed data with predicted values for the range of possible values of $f$, i.e. between 0.03 and 0.6, where 0.6 represents the free-streaming limit.

Three different laser intensity regimes are investigated ($2 \times 10^{13}$, $2 \times 10^{14}$, and $1 \times 10^{15}$ W/cm$^2$) using a trapezoidal (100 ps rise and fall times, 400 ps flat top) 0.532 μm laser pulse. Two laser absorption models are also compared: the traditional inverse bremsstrahlung (IB) absorption and an electromagnetic wave solver (EMS). The first calculates the local absorption factor, $\alpha$, as the laser light penetrates into the target, and requires a free parameter to start the simulation. This parameter legislates a fraction of the penetrating laser energy to be deposited at the critical density surface. The second method solves the time evolution of the Helmholtz equations for electromagnetic waves in an inhomogeneous dielectric. It is shown that the predicted shock speed is sensitive to $f$ in the range $0.03 \lesssim f \lesssim 0.08$ at $2 \times 10^{13}$ W/cm$^2$, and $0.03 \lesssim f \lesssim 0.15$ at $1 \times 10^{15}$ W/cm$^2$. A shock speed of $(3 \times 10^6 \pm 5\%)$ cm/s is predicted using the EMS method and laser irradiance of $(2 \times 10^{14} \pm 10\%)$ W/cm$^2$ with $f \simeq 0.090_{-0.022}^{+0.005}$. The IB method does not give a unique solution. The same irradiance gives the required speed with $f \simeq 0.045_{-0.007}^{+0.006}$ assuming 1% energy dump at the critical density, and with $f \simeq 0.03$ assuming 10% dump. For this reason, the EMS method is preferred.
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This thesis is dedicated to my grandmother, Virginia Smith, who died in 1994.
Chapter 1

Introduction

The development of high power pulsed lasers has provided the opportunity to produce and study high pressure shock waves. Complex computer codes (such as FILM [1], LAPLAS [2], LASNEX [3], MEDUSA [4], and SAGE [5]), which model the physical processes involved in laser-matter interactions, are needed not only to help interpret the results, but also to design new experiments. These programs incorporate existing theories of laser light absorption, mass ablation, thermal conduction, and material properties. However, the physics governing shock-compressed solids is still not well understood.

Shock-compressed solids are characterized by densities greater than solid density, high temperatures and pressures. They can be created as a result of the interaction of a high intensity, short pulse laser beam incident on a solid target. Four distinct regions in the target are observed during this process (Figure 1.1): the coronal plasma, ablation zone, shock compressed solid, and solid material. Laser radiation is absorbed at the cold target surface. The material is heated and ionized forming a dense plasma. This plasma is bounded by the solid target on one side and a vacuum on the other; therefore, the plasma expands rapidly into the vacuum as a rarefaction wave at the ion-acoustic speed. This underdense plasma region is called the corona. It is characterized by low density ($<10^{23}$ cm$^{-3}$) and high temperature ($>100$ eV). This is also referred to as the absorption region since absorption of laser light only occurs in the corona. EM radiation can penetrate only up to the critical density layer, $n_{cr}$, which is defined as the density where the plasma frequency equals the laser frequency. At this surface, the remaining
Figure 1.1: Schematic of a laser beam irradiating a solid target, showing the formation of a shock front. $n_{cr}$ is the critical density surface.
laser energy is reflected back out into the corona where further absorption can occur.

The next region of the target is the conduction or ablation zone. A large temperature difference now exists between the cold solid target and the critical density layer. Thus, the absorbed energy in the corona is conducted into the cold target, mainly by electron thermal transport, causing ablation of the solid up to the ablation front. This heated material expands out into the coronal plasma and into the vacuum.

To conserve the momentum of the outward flux of the ablated material, a shock wave develops and propagates through the target faster than the ablation front. Initially a compression wave is generated in the target. The speed of propagation of this wave is given by the local isentropic sound speed ($c_s$),

$$c_s = \sqrt{\left(\frac{\partial P}{\partial \rho}\right)_s} \quad (1.1)$$

where $P$ is the pressure which is a function of density, and $\rho$ is the material density. For adiabatic compression in a shock wave,

$$c_s = \sqrt{\frac{\gamma P}{\rho}} \quad (1.2)$$

where $\gamma$ is the specific heat. The tail end of the pressure disturbance is traveling through higher density material since the wave has already compressed this material. Thus the tail end travels faster than the wave front (Figure 1.2).\(^1\) This leads to steepening of the density gradient, and thus the temperature and pressure profiles, as the rear of the disturbance catches up to the leading edge of the wave front. Eventually a discontinuity is observed at the wave front which is termed a shock wave front. As the shock wave continues to propagate into the target it compresses more material.

\(^1\)Volume $V \propto \rho^{-1}$. Thus for an adiabatic process, $PV^\gamma=\text{constant}$ gives $P\rho^{-\gamma}=\text{constant}$. Therefore, from Equation (1.2), $c_s \propto \rho^{(\gamma-1)/2}$. For an ideal, monatomic gas $\gamma=5/3$, thus $c_s \propto \rho^{1/3}$. 
Figure 1.2: Schematic showing the steeping of a shock profile due to higher sound speed in more dense material.
1.1 Thesis Motivation

Thermal conduction in the ablation zone is modeled using the Spitzer-Härm diffusion law. This description was derived assuming small temperature gradients. However, steep temperature gradients (temperature scale lengths only a few electron mean free paths) are often observed near the critical density surface due to rapid expansion of the plasma into the vacuum. This results in the classical treatment overestimating the heat flux, $Q$. To prevent unrealistic predictions of heat flow, a free parameter called the flux limiter, $f$, was incorporated into the thermal conduction model as a simple solution to reduce $Q$:

$$Q = f n_e k_B T_e v_{te}$$  \hspace{1cm} (1.3)

where $n_e$ is the electron number density, $k_B$ is Boltzmann's constant, $T_e$ is the electron temperature, and $v_{te}$ is the electron thermal velocity. A value of $f=0.6$ corresponds to the free-streaming limit, the maximum heat flux that can be carried by electrons with a Maxwellian velocity distribution. A lower value of $f$ is then used to legislate heat flow to match experimental observations.

There is still some disagreement in the value of $f$ required to model heat conduction. Many authors [1, 2, 5, 6, 7, 8, 9] found that strong inhibition ($0.01 \lesssim f \lesssim 0.05$) was required to best fit their data, while others [10, 11, 12, 13, 14] maintained that less inhibition was needed ($0.1 \lesssim f \lesssim 0.3$). This discrepancy may be due to the differences in pulse rise time. Table 1.1 shows that the two regimes of flux inhibition, large values of $f$ versus small values, seem to correspond to nanosecond and picosecond pulses, respectively. The trapezoidal pulse used by Montgomery et al. is a nanosecond pulse in total duration but the rise time is only 100 ps. The trend is not surprising since heat inhibition occurs due to steep temperature gradients. With the more rapid heating due to a short pulse the gradients will build up more quickly as there is less time for heat conduction to dissipate the energy into the target. There is one exception to this given by the work of Town
et al. who compared a Fokker-Planck model with flux limited Spitzer heat flow using a much shorter pulse (500 fs) which indicated little flux inhibition ($f=0.3$) was needed.

Table 1.1: List of the estimate for flux limiter by different authors.

<table>
<thead>
<tr>
<th>Author</th>
<th>$f$</th>
<th>$\lambda$ (µm)</th>
<th>Irradiance (W/cm²)</th>
<th>Pulse Width</th>
<th>Power (GW)</th>
<th>Energy (J)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Fabbro et al. [1]</td>
<td>0.01-0.05</td>
<td>0.26</td>
<td>$1.5 \times 10^{14}$</td>
<td>60 ps</td>
<td>6</td>
<td>0.3</td>
</tr>
<tr>
<td></td>
<td>0.01-0.04</td>
<td>0.53</td>
<td>$5 \times 10^{14}$</td>
<td>80 ps</td>
<td>10</td>
<td>0.8</td>
</tr>
<tr>
<td></td>
<td>0.006-0.03</td>
<td>1.06</td>
<td>$1 \times 10^{15}$</td>
<td>100 ps</td>
<td>30</td>
<td>3</td>
</tr>
<tr>
<td>Amiranoff et al. [2]</td>
<td>0.02</td>
<td>1.3</td>
<td>$2 \times 10^{14}$</td>
<td>300 ps</td>
<td>300</td>
<td>100</td>
</tr>
<tr>
<td>Burgess et al. [5]</td>
<td>0.03</td>
<td>1.06</td>
<td>$\sim 10^{15}$</td>
<td>20 ps</td>
<td>75</td>
<td>1.5</td>
</tr>
<tr>
<td>Montgomery et al. [6]</td>
<td>0.03</td>
<td>0.35</td>
<td>$1 \times 10^{15}$</td>
<td>20 ns flat top 100 ps rise*</td>
<td>5.5</td>
<td>11,000</td>
</tr>
<tr>
<td>Benattar et al. [7]</td>
<td>$\lesssim 0.05$</td>
<td>1.06</td>
<td>$2 \times 10^{15}$</td>
<td>100 ps</td>
<td>100</td>
<td>1-8</td>
</tr>
<tr>
<td>Malone et al. [9]</td>
<td>0.03-0.1</td>
<td>1.06</td>
<td>$5 \times 10^{15}$</td>
<td>25 ps</td>
<td>160-1600</td>
<td>4-40</td>
</tr>
<tr>
<td>Wyndham et al. [10]</td>
<td>0.15±0.05</td>
<td>10.6</td>
<td>$\leq 1.3 \times 10^{11}$</td>
<td>75 ns</td>
<td>$\leq 0.2$</td>
<td>15</td>
</tr>
<tr>
<td>Goldsack et al. [11]</td>
<td>0.1</td>
<td>0.53</td>
<td>$2 \times 10^{12}-9 \times 10^{13}$</td>
<td>1.0 ns</td>
<td>40</td>
<td>40</td>
</tr>
<tr>
<td>Town et al. [12]</td>
<td>0.3</td>
<td>0.25</td>
<td>$\sim 10^{15}$</td>
<td>0.5 ps</td>
<td>not given</td>
<td>not given</td>
</tr>
</tbody>
</table>

*All pulses are Gaussian (with the given full width half maximum), except the pulse indicated which is trapezoidal.

Experimental evidence for heat flux inhibition has been found through comparison of mass ablation rates [1, 9, 11], x-ray emission [6, 9], electron density [5, 7, 10], and shock front trajectory [2] measurements with numerical simulations. Refer to Table 1.1 for the relevant laser parameters and value of $f$ which provided the best match between experimental data and numerical simulations.

Early heat transport measurements, such as those performed by Malone et al. [9], were
time-integrated, but laser-produced plasmas are time and space-dependent. Use of x-ray streak spectroscopy and Faraday cups (which measure ion current) by Goldsack et al. [11] provided temporal resolution for mass ablation. However, they used layered targets, glass spheres and shells coated with plastic and aluminum, so their results may be affected by beam inhomogeneities and target imperfections. Mass ablation measurements made by Fabbro et al. [1] may be affected by hot spots in the laser beam and nonuniformities in the target foil since they are measuring transmitted light through very thin polystyrene foils (0.04-10 μm).

Montgomery et al. [6] observed the x-ray ring emission from a titanium layer on Formvar foil as a function of time after laser beam burnthrough. This was the first direct measurement of a radially propagating heat front. Though the simulations and experimental data were in qualitative agreement, a later burnthrough time was predicted in the calculations.

Interferometry techniques were used to measure electron density profiles by Burgess et al. [5] and Benattar et al. [7]. The density profile of Burgess et al. shows reasonable agreement with simulations only if the probe pulse actually arrived ~20 ps later than experimentally determined. They suggest that this may be due to inadequate modeling of the hydrodynamics at early times. Also, the simulations do not show the observed density plateau around $n_{cr}/4$. A plot of the density trajectory for $n_{cr}$ and $n_{cr}/10$ was provided as well. Poor agreement was found except at early times. The calculations all used the ideal gas equation of state for carbon, which is inaccurate for laser-produced plasmas. One calculation was shown using SESAME data tables for the equation of state, but this calculation was only shown for the case of classical heat flux, and this also did not agree well with the experimental data. Density and temperature measurements of a plasma, created using slow Z-pinch in hydrogen, were made by Wyndham et al. [10] using Thompson scattering and was one of the first to provide experimental evidence for
larger heat flux.

Amiranoff et al. [2] reported an experiment observing the shock front trajectory produced in Plexiglass using two different laser intensities. They compared numerical predictions using $f=0.02$, and 0.6 with experimental data. A one-dimensional code was used for the lower irradiance, and a two-dimensional code the the higher one. The flux limited results ($f=0.02$) gave a reasonable match with the experimental observations. However, electrical breakdown [15] of the target was observed. Electrical breakdown occurs because at high intensities electric fields are created in the target and on the surface which can exceed the breakdown strength of insulating material. This effect obscures (an opaque cloud is observed in the target) the shock front at early times permanently damaging the target, there is an increase in the energy dissipation along the breakdown channels, and the homogeneity of the material is destroyed. It is not apparent that this phenomenon was included in the simulations by Amiranoff et al. Thus the validity of their model calculation is in serious question.

Theoretical attempts to determine the optimal value for the flux limiter have also been made through comparison of the Spitzer-Härm heat conduction model (which was used in the previous examples) to alternate methods of simulating laser-matter interactions. One such technique numerically solves the electron kinetic equation: the Fokker-Planck equation. The results from Town et al. [12], Matte and Virmont [13], and Bell et al. [14] provide evidence of reduction in thermal conduction from the classical Spitzer-Härm flux by a factor of $f=0.3$, 0.3, and 0.1, respectively. Mason [8] used a self-consistent Monte Carlo model for electron thermal transport, and found that greater inhibition was required. Some agreement between the two models was found with $f=0.03$. Experimental data was not provided by these authors to verify the validity of the models.

This thesis will explore the feasibility of using shock speed as a probe of flux-limited thermal conduction. Though this is not a direct measurement of thermal transport it
has the advantage of diagnostic simplicity. X-ray diagnostics, ion calorimeters, streak cameras, and layered targets are not required. Shock speed may easily be determined from the shock transit time in a simple, bulk target (aluminum for this work). This method is also less prone to non-uniformities in the laser beam.
Chapter 2

Physical Processes in Laser-Matter Interactions

2.1 Laser Absorption Mechanisms

2.1.1 Inverse Bremsstrahlung Absorption

Inverse bremsstrahlung (IB) absorption, or collisional absorption, occurs when electrons oscillate in the laser’s electric field and collide with ions (thus it is a free-free interaction). The ions are required in this process to conserve momentum. The coherent energy of oscillation of the electrons is then converted into random thermal energy. Absorption occurs by IB only in the underdense plasma up to the critical density surface, where resonance absorption may be observed. The remaining energy is reflected and can be absorbed via IB on the way out of the target.

To understand why laser radiation is deposited only up to the critical density surface, consider the plane wave equation for the laser electric field, $E_L$:

$$E_L = E_{L0} \exp(i(k_L \cdot r - \omega_L t))$$  \hspace{1cm} (2.1)

where the propagation of the laser electric field in a plasma is given by the dispersion relation

$$\frac{k_L^2 c^2}{\omega_L^2} = \epsilon = 1 - \frac{\omega_p^2}{\omega_L^2 + \nu_{ei}^2} + i \frac{\nu_{ei} \omega_p^2}{\omega_L (\omega_L^2 + \nu_{ei}^2)}$$  \hspace{1cm} (2.2)

where $k_L$ is the wave number, $\omega_L$ is the laser frequency, $\epsilon$ is the plasma dielectric constant, $c$ is the speed of light, $\nu_{ei}$ is the electron-ion collision frequency, and $\omega_p$ is the plasma
frequency given by
\[ \omega_p = \sqrt{\frac{n_e e^2}{m_e \epsilon_0}} \]  

(2.3)

where \( e \) is the charge on an electron, \( n_e \) is the electron number density, \( m_e \) is the mass of an electron, and \( \epsilon_0 \) is the permittivity of free space. At optical frequencies, usually \( \nu_{ei} \ll \omega_L \) [16, 17, 18] with \( \nu_{ei} \approx 10^{14} \text{ s}^{-1} \). In this thesis, \( \omega_L = 3.5 \times 10^{15} \text{ rad/s} \) and \( \nu_{ei} \approx 1 \times 10^{14} \text{ s}^{-1} \). Thus, Equation (2.2) may be written, to the first approximation, as

\[ \frac{k_L^2 c^2}{\omega_L^2} = \epsilon \approx 1 - \frac{\omega_p^2}{\omega_L^2} \]  

(2.4)

Equation (2.4) shows that as the EM wave propagates from lower to higher densities the wave number becomes smaller. At the critical density, \( n_{cr} \), the natural frequency of oscillation of electrons in a plasma, \( \omega_p \), is the same as the laser frequency, \( \omega_L \). Thus the real part of \( k_L \) goes to zero and the wave cannot propagate further. For electron densities greater than \( n_{cr} \), \( \omega_p > \omega_L \), and \( k_L \) is purely imaginary. Thus the electric field is an evanescent wave. From Equation (2.3), \( n_{cr} \) (in \( \text{cm}^{-3} \)) can be found,

\[ n_{cr} = \frac{m_e \epsilon_0 \omega_L^2}{e^2} = \frac{1.1 \times 10^{21}}{\lambda_L^2} \]  

(2.5)

with the laser wavelength, \( \lambda_L \), given in \( \mu\text{m} \).

**Traditional Treatment**

Two techniques for modeling laser light absorption will be examined: the traditional method, and using an EM wave solver. This section will describe the first method, and the next section will describe the second.

Initially the target is at solid density which is greater than the critical density. Therefore, to start the absorption process the traditional method requires a free parameter called the fractional energy dump. This parameter calculates a user-legislated fraction
of the laser energy that reaches the critical density layer to be absorbed at that layer. In the initial stage, this is the energy that is absorbed at the solid target surface.

The traditional method determines the local energy absorption in the target. The spatial attenuation of the incident EM wave is described by the absorption coefficient $\alpha$ [4]:

$$\alpha = \frac{13.51 \beta^2}{\lambda_L^2 (1 - \beta)^{1/2}} \frac{Z \ln \Lambda}{T_e^{3/2}}$$

where $\lambda_L$ is given in m, $Z$ is the plasma ionization state, $T_e$ is the electron temperature given in K, and

$$\beta = \frac{n_e}{n_{cr}} < 1$$

with

$$n_e = Z n_i$$

where the ion number densities $n_i$ (in m$^{-3}$) is given by

$$n_i = \frac{\rho}{A m_p}$$

where $\rho$ is the plasma density, $A$ is the average atomic mass number of the target material, and $m_p$ is the proton mass, and $\ln \Lambda$ is the Coulomb logarithm [19]:

$$\ln \Lambda = 16.34 - \ln(T_e^{3/2} n_e^{-1/2} Z^{-1})$$

From Equation (2.6) it is apparent that IB absorption is strongest for high density, low temperature, high $Z$ plasmas. In this regime, most of the laser energy is absorbed before reaching the critical density. The absorbed power is given by

$$A_{abs}(r, t) = \frac{1}{\rho} \frac{\partial \Phi_L}{\partial r}$$

where the local laser intensity $\Phi_L$ reaching position $r$ from the plasma boundary $R_0$ is:

$$\Phi_L(r, t) = \Phi_L(R_0, t)[1 - \exp \int_r^{R_0} -\alpha(r', t)dr']$$
EM Wave Solver

This method for describing the laser light deposition solves the Helmholtz equations for EM waves [16] in a plasma [20], treating the target material as an inhomogeneous dielectric with the complex dielectric function [17]:

\[
e(\omega) = 1 + i \frac{4\pi \sigma(\omega)}{\omega}
\]  

(2.13)

where \( \sigma(\omega) \) is the electrical conductivity at the frequency \( \omega \). The boundary conditions for incident and reflected waves in a vacuum, and an evanescent wave in the target are used to numerically solve the Helmoltz equations at each time step of the calculation, providing the complex electric field amplitude and reflectivity. Thus, unlike the previous method, the EM wave solver does not need to use a legislated energy dump to begin the absorption process as an evanescent wave penetrates into a cold solid. The energy deposition rate can then be determined

\[
\langle \mathbf{E}_L \cdot \mathbf{J} \rangle = \frac{1}{2} Re(\sigma)|\mathbf{E}_L|^2
\]

(2.14)

which is the heat source term \( (A_{\text{las}}) \) for the hydrodynamic equations. The Drude approximation is used to determine the electrical conductivity \( \sigma(\omega) \) [21],

\[
\sigma(\omega) = \frac{\omega_p^2}{4\pi(\nu_{\text{ei}} - i\omega)}
\]

(2.15)

where the plasma frequency \( \omega_p \) is given by

\[
\omega_p^2 = 4\pi Z n_i e^2 / m
\]

(2.16)

and the electron-ion collision rate \( \nu_{\text{ei}} \) is

\[
\nu_{\text{ei}} = Z n_i e^2 / m \sigma_0
\]

(2.17)

where \( \sigma_0 \) is the dc conductivity determined using values predicted by Lee and More [22] or Perrot and Dharma-wardana [23] or Rinker [24].
This local description of the dielectric function works well for s-polarized light. However, for p-polarized light resonance absorption is observed and energy deposition at $\nu_{cr}$ is limited by the density gradient scale length. The density gradient steepens with the rapid expansion of the plasma. Thus the regime where resonance absorption occurs narrows and this spatial length is ultimately determined by the mesh size chosen for the calculation. Use of a nonlocal dielectric function $\epsilon^*(z)$ avoids this unphysical limit resulting from the mesh size by spatially averaging over a minimum characteristic scale length $\lambda_s$. The scale length is chosen to be the greater of the degeneracy-corrected Debye length $\lambda_D$ or the interatomic distance $r_0$. This nonlocal dielectric function is found using

$$\epsilon^*(z) = \frac{1}{\lambda_s} \int_{-\lambda_s/2}^{\lambda_s/2} \epsilon(z + z')dz'$$

(2.18)

**Langdon Factor**

At high laser intensities the electron velocity distribution is no longer Maxwellian since the electrons absorb the laser energy at a rate greater than their relaxation towards the Maxwellian distribution, i.e.

$$\nu_{ei}v_{osc}^2 > \nu_{ee}v_{te}^2 \leftrightarrow Z\frac{eE_L}{m_e\omega_Lv_{te}} > 1$$

(2.19)

where $v_{osc}$ is the electron peak velocity of oscillation in the electric field, $\nu_{ee} (\sim \nu_{ei}/Z)$ is the electron-electron collision frequency, and $v_{te}$ is the electron thermal velocity. Those electrons with $\nu_{ei}$ close to the laser frequency have the largest contribution towards absorption, and since $\nu_{ei}(v) < \omega_L$ and $\nu_{ei}(v) \propto v^{-3}$ these electrons have a slower velocity than $v_{te}$. However, for high laser intensities when the above is satisfied, there are relatively fewer low velocity electrons and the electron distribution is distorted. This results in reduced laser energy absorption up to a factor of 2. This nonlinear effect can be ignored if

$$\frac{v_{osc}^2}{v_{te}^2} = 4 \times 10^{-6} \frac{I_L\lambda_L^2}{T_e} << 1$$

(2.20)
where $I_L$ is the laser intensity (in W/m$^2$), $\lambda_L$ is the laser wavelength (in \(\mu\)m), $T_e$ is the electron temperature (in keV). For high irradiances when it cannot be neglected the Langdon opacity correction factor [25] has been included in the EMS calculation such that

$$\sigma'_o = \sigma_o / f_{\text{lang}}$$

where

$$f_{\text{lang}} = 1 - \frac{0.553}{(1 + \frac{0.27}{\nu_{\text{esc}}/\nu_e})^{3/4}}$$

2.1.2 Resonance Absorption

Resonance absorption occurs at the critical density surface when the electric field of the incident laser resonantly excites the electron plasma waves (Langmuir waves). Thus this absorption process dominates in the regime where IB absorption is inefficient and most of the laser energy reaches the critical density layer, i.e. for high laser intensities, long laser wavelengths, and short plasma scale lengths.

To understand this process, consider p-polarized light. $E_L$ is parallel to the plane of incidence, obliquely incident on a plasma with density gradient $\nabla n_e$. The component of the laser electric field parallel to the density gradient can penetrate through the critical density surface, creating a charge separation that drives a charge density fluctuation $\delta n_e$ at the laser frequency $\omega_L$. At the critical density $\omega_L = \omega_p$ and the electron plasma waves are resonantly excited. This can be illustrated using Maxwell’s equation

$$\nabla \cdot [\epsilon(x)E_L] = 0$$

where $\epsilon(x)$ is found from the dispersion equation

$$\epsilon(x) = 1 - \frac{\omega_p^2(x)}{\omega_L^2}$$
Expanding Equation (2.23) gives

\[ \varepsilon(x) \nabla \cdot \mathbf{E}_L + (\mathbf{E}_L \cdot \nabla)\varepsilon(x) = 0 \]  

(2.25)

Now \( \nabla \varepsilon(x) \propto \nabla n_e \), thus by comparison with Poisson’s equation

\[ \nabla \cdot \mathbf{E}_L = -\frac{e\delta n_e}{\varepsilon_o} \]  

(2.26)

it can be seen that

\[ \frac{e\delta n_e}{\varepsilon_o} \propto \frac{\mathbf{E}_L \cdot \nabla n_e}{\varepsilon(x)} \]  

(2.27)

Thus a density perturbation arises when \( \varepsilon(x) \approx 0 \) or \( \omega_p = \omega_L \), and when \( \mathbf{E}_L \) is in the same direction as the density gradient.

For s-polarized light, the laser electric field vector is perpendicular to the plane of incidence; therefore, there is not a component parallel to the density gradient and resonant absorption does not occur.

In numerical simulations, when the traditional method is used, this absorption method is often modeled as a constant fraction of laser energy being deposited at the critical density layer [26]. Thus

\[ A_{tas} = a\Phi_L(R_0, t) \]  

(2.28)

where \( a \) is the fractional energy dump. The EM wave solver is able to calculate resonance absorption more rigorously. The laser light used for this thesis is p-polarized. However it is at normal incidence; therefore, resonance absorption is not observed.

### 2.1.3 Anomalous Absorption

As illustrated with resonance absorption, laser-driven fluctuations in plasma density excite the natural modes of oscillation of the plasma. Either electron plasma waves or ion-acoustic waves may be resonantly excited, resulting in absorption or scattering of the laser
energy by the plasma. Four such mechanisms will be described here: parametric-decay instability, two-plasmon decay instability, stimulated Raman scattering, and stimulated Brillouin scattering. [18]

Parametric-decay instability is the decay of the incident wave (with frequency $\omega_0$ and wave number $k_0$) into an electron plasma wave ($\omega_{ek}$, $k_{ek}$) and an ion-acoustic wave ($\omega_{ia}$, $k_{ia}$). Fluctuations in ion density couples energy into an electron plasma wave. The electron plasma wave then oscillates at the same frequency as the laser light resulting in spatial variations in the electric field intensity. This gradient leads to a ponderomotive force which enhances the ion density fluctuations, completing the feedback loop. Frequency and wave number matching conditions must be met for strong coupling of the laser light into the excited wave, allowing the instability to grow:

\[
\begin{align*}
\omega_0 &\to \omega_{ek} + \omega_{ia} \\
k_0 &\to k_{ek} + k_{ia}
\end{align*}
\]  

(2.29)

The minimum frequency of a light wave in a plasma is $\omega_{pe}$ (electron plasma frequency) which is proportional to $n_e^{-1/2}$, and $\omega_0 \propto n_{cr}^{-1/2}$. Also, the frequency matching condition is satisfied for $n_e \approx n_{cr}$.

Two-plasmon decay instability results when the laser light resonantly decays into two electron plasma waves (plasmons). The frequency and wave number matching conditions now are:

\[
\begin{align*}
\omega_0 &\to \omega_{ek1} + \omega_{ek2} \\
k_0 &\to k_{ek1} + k_{ek2}
\end{align*}
\]  

(2.30)

Since the frequencies of the two plasmons are $\sim \omega_{pe}$ (electron plasma frequency), the frequency matching condition requires that this instability occurs at $n \approx n_{cr}/4$. During both parametric-decay instability and two-plasmon decay instability, energy is transferred to plasma waves, representing absorption by the plasma.
Stimulated Raman scattering involves the resonant decay of the incident light wave into a scattered light wave and an electron plasma wave. Instability results when a feedback loop arises, resulting in the growth of these waves. To understand this process, consider a small density fluctuation $\delta n$, associated with an electron plasma wave, which is in the direction of propagation of the laser electric field $E_L$. The oscillations of the electrons in $E_L$ generate a transverse current $\delta J = -e^2 E_L \delta n / m \omega_0$. If the matching conditions,

$$\omega_0 \rightarrow \omega_s + \omega_{ek}$$

$$k_0 \rightarrow k_s + k_{ek}$$

are met, then this current generates a scattered light wave with electric field $\delta E$. This wave then interferes with the incident wave producing a change in wave pressure, $\nabla(E^2 / 8\pi) = \nabla(E_L \cdot \delta E) / 4\pi$, that enhances the density fluctuations. This feedback loop results in increased scattering and energy absorption by the plasma. The minimum frequency of a light wave in a plasma is $\omega_{pe}$ (electron plasma frequency); therefore, according to the frequency matching condition, this instability requires $\omega_0 \sim 2\omega_{pe}$ or $n \lesssim n_{cr}/4$.

Stimulated Brillouin scattering is similar to stimulated Raman scattering, except that the incident wave is coupled into a scattered wave and an ion-acoustic wave, rather than an electron plasma wave. The frequency and wave number matching conditions for this process are:

$$\omega_0 \rightarrow \omega_s + \omega_{ia}$$

$$k_0 \rightarrow k_s + k_{ia}$$

An ion-acoustic wave oscillates with a much lower frequency than $\omega_0$; therefore, the instability can be observed throughout the underdense plasma, i.e. $n < n_{cr}$.

As shown, the frequency or wave number matching conditions determine the region of the plasma where the instability may occur. There is also a minimum incident laser
intensity required such that the excited waves grow. This pumping energy is greater than the damping of these waves. The threshold intensities are provided in Table 2.1 as functions of the plasma parameters $T_e$ (in eV), $n_e$ (in cm$^{-3}$), $\lambda_L$ (in $\mu$m), and $L=c_s t$ (the density gradient scale length, in $\mu$m). Three intensity regimes will be investigated in this work: $2 \times 10^{13}$, $2 \times 10^{14}$, and $1 \times 10^{15}$ W/cm$^2$. Table 2.2 illustrates that the parametric-decay instability may play a role for all irradiances, however since it only occurs near critical density this will only be important when inverse bremsstrahlung absorption is not significant. The 2-plasmon decay instability or stimulated Brillouin scattering may affect results for the larger two irradiances. Stimulated Raman scattering is not expected to be significant.

Table 2.1: Partial list of instabilities which can be driven by laser light propagating through a plasma, resulting in absorption or scattering of light.

<table>
<thead>
<tr>
<th>Name</th>
<th>Process</th>
<th>Density Region (cm$^{-3}$)</th>
<th>Threshold Intensity (W/cm$^2$) [18, 27]</th>
</tr>
</thead>
<tbody>
<tr>
<td>parametric-decay instability (PDI)</td>
<td>$\omega_0 \to \omega_{ck} + \omega_{ia}$</td>
<td>$n_e \sim n_{cr}$</td>
<td>$8.45 \times 10^{11} T_e L^{-1}$</td>
</tr>
<tr>
<td>2-plasmon decay instability (2PDI)</td>
<td>$\omega_0 \to \omega_{ck1} + \omega_{ck2}$</td>
<td>$n_e \sim n_{cr}/4$</td>
<td>$3 \times 10^{13} Z T_e^{-1/2} \lambda_L^{-3}$</td>
</tr>
<tr>
<td>stimulated Raman scattering (SRS)</td>
<td>$\omega_0 \to \omega_s + \omega_{ck}$</td>
<td>$n_e &lt; n_{cr}/4$</td>
<td>$5 \times 10^{17} L^{-4/3} \lambda_L^{-2/3}$</td>
</tr>
<tr>
<td>stimulated Brillouin scattering (SBS)</td>
<td>$\omega_0 \to \omega_s + \omega_{ia}$</td>
<td>$n_e &lt; n_{cr}$</td>
<td>$7.5 \times 10^{12} T_e \lambda_L^{-1} L^{-1} \left( \frac{n_{ee}}{n_e} \right)$</td>
</tr>
</tbody>
</table>

2.2 Electron Thermal Transport

2.2.1 Spitzer-Härm Theory of Thermal Conductivity

Laser energy is absorbed in plasma corona exciting the plasma electrons. The heated electrons then transport the energy from this absorption region towards the cold solid
Table 2.2: Typical plasma conditions predicted by the EMS and the corresponding threshold intensities for various anomalous absorption processes. \( n_{cr} = 3.89 \times 10^{21} \text{ cm}^{-3} \) for \( \lambda = 0.532 \mu\text{m} \).

<table>
<thead>
<tr>
<th>Laser Intensity (W/cm(^2))</th>
<th>Density Region</th>
<th>( T_e ) (eV)</th>
<th>( Z )</th>
<th>( L ) (μm)</th>
<th>Instability</th>
<th>Threshold Intensity (W/cm(^2))</th>
</tr>
</thead>
<tbody>
<tr>
<td>( 2 \times 10^{13} )</td>
<td>( n_e \sim n_{cr} )</td>
<td>225</td>
<td>12</td>
<td>30</td>
<td>PDI</td>
<td>6.3 \times 10^{12}</td>
</tr>
<tr>
<td></td>
<td>( n_e \sim n_{cr}/4 )</td>
<td>290</td>
<td>12.5</td>
<td>34</td>
<td>2PDI</td>
<td>1.5 \times 10^{14}</td>
</tr>
<tr>
<td></td>
<td>( n_e &lt; n_{cr}/4 )</td>
<td>290</td>
<td>12.5</td>
<td>34</td>
<td>SRS</td>
<td>6.9 \times 10^{15}</td>
</tr>
<tr>
<td>( 2 \times 10^{14} )</td>
<td>( n_e \sim n_{cr} )</td>
<td>530</td>
<td>13</td>
<td>46</td>
<td>PDI</td>
<td>9.7 \times 10^{12}</td>
</tr>
<tr>
<td></td>
<td>( n_e \sim n_{cr}/4 )</td>
<td>766</td>
<td>13</td>
<td>56</td>
<td>2PDI</td>
<td>9.4 \times 10^{13}</td>
</tr>
<tr>
<td></td>
<td>( n_e &lt; n_{cr}/4 )</td>
<td>766</td>
<td>13</td>
<td>56</td>
<td>SRS</td>
<td>3.6 \times 10^{15}</td>
</tr>
<tr>
<td></td>
<td>( n_e &lt; n_{cr} )</td>
<td>530</td>
<td>13</td>
<td>46</td>
<td>SBS</td>
<td>1.6 \times 10^{14}</td>
</tr>
<tr>
<td>( 1 \times 10^{15} )</td>
<td>( n_e \sim n_{cr} )</td>
<td>1000</td>
<td>13</td>
<td>66</td>
<td>PDI</td>
<td>1.3 \times 10^{13}</td>
</tr>
<tr>
<td></td>
<td>( n_e \sim n_{cr}/4 )</td>
<td>1500</td>
<td>13</td>
<td>78</td>
<td>2PDI</td>
<td>6.7 \times 10^{13}</td>
</tr>
<tr>
<td></td>
<td>( n_e &lt; n_{cr}/4 )</td>
<td>1500</td>
<td>13</td>
<td>78</td>
<td>SRS</td>
<td>2.3 \times 10^{15}</td>
</tr>
<tr>
<td></td>
<td>( n_e &lt; n_{cr} )</td>
<td>1000</td>
<td>13</td>
<td>66</td>
<td>SBS</td>
<td>2.1 \times 10^{14}</td>
</tr>
</tbody>
</table>

in an attempt to equalize the temperature. Due to collisions with other electrons, the electrons are not all traveling with the same velocity, but have a Maxwellian distribution [28]:

\[
f(v) = \frac{n_e}{(2\pi)^{3/2}v_{te}^3} e^{-m_e v^2/2v_{te}^2}
\]  

where \( v_{te} = \sqrt{k_B T_e/m_e} \), \( v^2 \equiv (v_x^2 + v_y^2 + v_z^2) \). The electron density at time \( t \) and position \( r \) by definition is given by

\[
n_e(r, t) = \int f(r, v, t) dv
\]  

where \( f(r, v, t) \) is the velocity distribution function, and \( dv \) is a three-dimensional volume element in velocity space. The heat flux, \( Q \), carried by the electrons is then:

\[
Q = \int f(r, v, t) \frac{m_e v^2}{2} dv
\]
The electron velocity distribution function $f(v)$ is described by the Boltzmann equation for electrons in the absence of magnetic fields [28]:

$$\frac{\partial f}{\partial t} + v \cdot \frac{\partial f}{\partial r} - \frac{eE}{m_e} \cdot \frac{\partial f}{\partial v} = \left( \frac{\partial f}{\partial t} \right)_{\text{coll}}$$

(2.36)

where $(\frac{\partial f}{\partial t})_{\text{coll}}$ is a Fokker-Planck collision term [29] representing electron-electron and electron-ion interactions. Kruer [18] provides a derivation for the electron-ion contribution to the collision term:

$$C_{ei} = A \frac{\partial}{\partial v} \left[ \frac{v^2 I - vv}{v^3} \cdot \frac{\partial f}{\partial v} \right]$$

(2.37)

where $A = (2\pi n Z e^4/m^2) \ln \Lambda$, and $I$ is the identity matrix. For high $Z$ plasmas, the more complex electron-electron collision term, $C_{ee}$, is negligible.

In the absence of any fields, the solution to Equation (2.36) is the Maxwellian distribution (Equation (2.33)). However, a self-consistent electric field, $E$, is produced due to the motion of the charged particles in the temperature gradient, and is related to the current density, $j$, by:

$$j = \sigma E + \alpha \nabla T = 0$$

(2.38)

where $\sigma$ and $\alpha$ are the appropriate coefficients given by Spitzer and H"{a}rm [29]. The net current must be zero or an electrostatic field would rise without limit. To compensate for the current due to electron thermal transport, a secondary electric field builds up producing a return current.

A small temperature gradient near $n_{cr}$ results in a small perturbation of the distribution function from the Maxwellian [30]:

$$f(v) = f_0(v) + f_1(v) \quad \text{with} \quad f_1 << f_0$$

(2.39)

where $f_0(v)$ is the unperturbed Maxwellian distribution, and $f_1(v)$ is the first order perturbation involved in heat flow. Higher order perturbations are ignored. The solution
Chapter 2. Physical Processes in Laser-Matter Interactions

(Equation (2.39)) to the second order differential equation for \( f_1(v) \) (Equation (2.36)) is substituted into Equation (2.35) and upon comparison with the classical form of heat flux (Equation (2.47)), the expression for the thermal conductivity, \( \kappa \), can be found. Using this method, the Spitzer-Harm conductivity, neglecting magnetic fields, was found to be [31] (in SI units)

\[
\kappa = 2 \times 10^5 \left( \frac{2}{\pi} \right)^{3/2} \frac{e^4 k_B (k_B T_e)^{3/2}}{m_e^{1/2} \epsilon^4 Z \ln \Lambda} \left[ \frac{0.095(Z + 0.24)}{1 + 0.24Z} \right]
\]  

(2.40)

where \( c \) is the speed of light.

To determine the form of \( f_1(v) \), the distribution function (Equation (2.39)) may be written as the first two terms of a Legendre polynomial [18, 30, 32, 33]:

\[
f(v) = f_0(v) + f_1(v) \cos \theta
\]

(2.41)

where \( \theta \) is the angle between the direction of electron motion (\( v \)) and heat flow (i.e. in the direction of \( \nabla T_e \) and \( E \)). Substitution of \( f(v) \) into Equation (2.36) and collecting only the terms proportional to (\( \cos \theta \)) provides a differential expression for \( f_1(v) \):

\[
\frac{\partial f_1}{\partial t} + v \frac{\partial f_0}{\partial z} - \frac{eE}{m} \frac{\partial f_0}{\partial v} = - \frac{2A}{v^3} f_1
\]

(2.42)

assuming, for simplicity, high \( Z \) plasma. Assuming steady-state conditions gives:

\[
f_1 = \frac{v^4}{2A} \left( \frac{eE}{mv} \frac{\partial f_0}{\partial v} - \frac{\partial f_0}{\partial z} \right)
\]

(2.43)

The electric field is obtained using this equation and imposing the condition of charge neutrality, i.e. \( J = (4\pi e/A) \int_0^\infty v^3 f_1 dv = 0 \). \( f_1 \) is then found to be [18]

\[
f_1 = f_0 \frac{v^4}{4mA} \left[ \frac{8}{v_e^2(z)} - \frac{v^2}{v_e^4(z)} \right] \frac{\partial T_e}{\partial z}
\]

(2.44)

or in terms of the perturbation parameter, \( \lambda_e/L_T \), where \( \lambda_e \) is the electron mean free path which increases with the electron velocity [18, 28]

\[
\lambda_e = v_{te}/\nu_{ei} \simeq (5.9 \times 10^7 T_e^{1/2}) \left( \frac{T_e^{3/2}}{3 \times 10^{-6} \ln \Lambda n_e Z} \right) \simeq \frac{2 \times 10^{13} T_e^2}{Z n_e \ln \Lambda} \quad \text{(cm)}
\]

(2.45)
with $T_e$ in eV, and $L_T = |T_e/\nabla T_e|$ is the temperature gradient scale length

$$f_1 = \frac{3 \times 10^{-6} m_e \lambda_e}{8 \pi e^4} \frac{T_e}{L_T} \left( \frac{v}{v_{te}} \right)^4 \left[ \left( \frac{v}{v_{te}} \right)^2 - 8 \right] \quad (2.46)$$

This equation shows that the perturbation term, $f_1$, depends strongly on the ratio of electron velocity to electron thermal velocity ($f_1 \propto (v/v_{te})^6$ when this ratio is large). At some velocity, depending on $\lambda_e/L_T$, $|f_1|$ can become greater than $f_0$. When $|f_1| > f_0$ it is no longer valid to truncate the expansion of $f$ to only the first two terms. Higher order terms are not negligible. For steeper temperature gradients (larger values of $\lambda_e/L_T$) this occurs at slower velocities. Shvarts et al. [33] show that for $\lambda_e/L_T = 0.002$ $f_1$ becomes greater than $f_0$ for $v \simeq 3v_{te}$ while for $\lambda_e/L_T = 0.1$ the theory breaks down at $v \simeq 2v_{te}$. Also, it is possible for $(\cos \theta)$ to be negative, giving a negative value of $f$ for $|f_1| > f_0$ which is clearly unphysical, thus emphasizing the need to include higher order terms in the expansion. Gray and Kilkenny [30] found that heat conduction results mainly from the motion of electrons with velocity in the range 2.3 to 3 times the thermal velocity, and $f(v)$ became negative in this range if $\lambda_e/L_T > 0.015$.

### 2.2.2 Flux-limited Thermal Conductivity

Classical heat flow is given by:

$$Q_{SH} = -\kappa \nabla T \quad (2.47)$$

with $\kappa$ being the Spitzer-Härm conductivity. The heat flow term in the velocity distribution function may be written in terms of heat flow [18]:

$$f_1 = f_0 \frac{Q_{SH}}{32 \pi n T_e v_{te}} \left( \frac{v}{v_{te}} \right)^4 \left[ \left( \frac{v}{v_{te}} \right)^2 - 8 \right] \quad (2.48)$$

To compensate for the breakdown of the theory used to derive $\kappa$, a phenomenological parameter called the flux limiter, $f$, is incorporated into the heat flow model to prevent
unrealistic heat flow, setting an upper limit on heat flux:

\[ Q' = \min[Q_{SH}, Q] \] (2.49)

where \( Q \) is given in Equation (1.3). A second option often used when modeling such inhibition in heat flux is to use the harmonic mean of the flux:

\[ Q' = (Q_{SH}^{-1} + Q^{-1})^{-1} \] (2.50)

Kruer [18] and the references he provides, suggests that the need for a flux limiter in the traditional Spitzer-Härm diffusion model of thermal transport is oversimplified by using only a single flux limiter. This description treats an intrinsically multigroup phenomenon as a single group. The use of only one group averages over the distribution of velocities, but the electron mean free path is energy-dependent. Thus electrons should be divided into energy groups, each assigned a diffusion coefficient and coupled together by imposing charge neutrality. The difficulty is how to incorporate this idea into existing hydrodynamical codes. An example of the need to divide the modeling into different regimes with different flux limiters is shown by Mason [8] who compares a self-consistent Monte Carlo model with a flux inhibited Spitzer-Härm diffusion model. He found that the coronal density profiles matched best when \( f = 0.1 \) or 0.15 was used, and the coronal temperatures agreed using \( f = 0.6 \). However, \( f \approx 0.03 \) was required to match the temperatures near the critical density. This suggests that separate flux limiters are needed for the internal and coronal regions.

Two physical explanations have been proposed for the reduction in heat flux at sharp temperature gradients: ion-acoustic turbulence [9, 30, 32, 34, 35] and self-generated magnetic fields [36, 37, 38, 39].
Ion-Acoustic Instability

Ion-acoustic instabilities modify heat flow by skewing the electron velocity distribution function from a Maxwellian. This can be seen in Figure 2.1 which illustrates the reduced distribution $F(v_{\|}) = \int f(v)dv_{\perp}$ ($v_{\|}$ is the velocity in the direction of heat flow) for electrons and ions. The ionic distribution is much narrower because the ions are taken to be considerably colder than the electrons as this is the regime when ion-acoustic instability may be observed. Heat flux is mainly carried by high energy electrons. The non-Maxwellian tail in the electron distribution for $v < 0$ is caused by low energy electrons which provide a return current necessary to compensate for the flow of the higher energy electrons into the target, and prevent a net current flow. This return current causes a drift, $v_d$, in the distribution function [30]:

$$v_d \simeq 0.44 \frac{\lambda_e}{L_T} v_{te} \quad (2.51)$$

The ion-acoustic wave velocity ($v_{ia}$), or phase velocity, is given by [18]:

$$v_{ia} = \left( \frac{Zk_BT_e + 3k_BT_i}{m_i} \right)^{1/2} \quad (2.52)$$

where $T_i$ and $m_i$ are the ion temperature and mass, respectively. If $T_e >> T_i$, then $v_{ia} >> v_{ti}$ and the phase velocity lies in the tail of the ion velocity distribution. This means that there are insufficient ions for Landau damping of the ion waves. Landau damping (or growth) of a wave results from the exchange of energy between the wave and particles traveling at velocities near the phase velocity. In plasmas there are charged particles traveling slightly faster and slower than this velocity. With a Maxwellian distribution of velocities there are more slow moving particles, compared to the wave velocity, than fast so the wave loses more energy to the particles than it gains, and Landau damping occurs. However, when the drift velocity, $v_d$, of the electrons exceeds $v_{ia}$ there are more electrons with higher energy and Landau growth is greater than Landau damping driving unstable
Figure 2.1: Electron and ion reduced distributions in the presence of an electron heat flux. When $v_D > v_{i\alpha}$, ion-acoustic instability may be observed.
ion-acoustic waves. The threshold for this instability may be written as

$$v_d > v_{ia} \sim \left( \frac{Zk_B T_e}{m_i} \right)^{1/2}$$

or

$$\frac{\lambda_e}{L_T} > 2.3 \left( \frac{Zm_e}{m_i} \right)^{1/2} \simeq 0.05 Z^{1/2} \tag{2.54}$$

where $m_i$ is the ion mass, and $v_d$ is given in Equation (2.51). An alternate way to determine if the wave will be resonantly excited is if the slope of the electron velocity distribution is positive at the phase velocity then Landau growth may occur. This growth can be significant if $\frac{\partial f}{\partial v} >> 0$. As can be seen in Figure 2.1, the slope of the electron velocity distribution is positive when $v_d > v_{ia}$.

Ion turbulence increases the electron-ion collision frequency, $\nu_{ei}$. Heat flux is then reduced since thermal conductivity is inversely proportional to $\nu_{ei}$ ($\kappa \sim n_e v_e \lambda_e = n_e \frac{v_e^2}{\nu_{ei}}$). Campbell et al. [40] measured time resolved ion current and found that about 50% of the absorbed laser energy was located in only a few percent of the target mass at the surface. The large concentration of energy in this mass may lead to poor energy transfer to the rest of the target. When ion-acoustic instability was included in their simulations, they were able to reproduce the characteristics of the ion velocity distributions. There is much disagreement as to the significance of the effect of ion-acoustic turbulence on heat flux. Some authors find significant inhibition [14, 34, 40] while others suggest that there is little effect [8, 18]. In this thesis, $\lambda_e/L_T \simeq 0.002, 0.005, 0.08$ for laser irradiances of $2 \times 10^{13}$, $2 \times 10^{14}$, and $1 \times 10^{15}$ W/cm$^2$, respectively. Thus, ion acoustic turbulence may only be observed in the latter case.

**Magnetic Fields**

A second mechanism that may affect heat transport is due to large self-generated magnetic fields that have been observed in laser-generated plasmas. There are several source
terms for these magnetic fields [36], the simplest occurring when the density and temperature gradients are not parallel. This is easily shown by incorporating the Lorentz force:

\[ \mathbf{F} = q(\mathbf{E} + \mathbf{u} \times \mathbf{B}) \]  

(2.55)

into Faraday’s Law:

\[ \nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \]  

(2.56)

where \( q \) is the charge on the particle, \( \mathbf{u} \) is the plasma flow velocity, and \( \mathbf{B} \) is the magnetic field. From the force exerted on electrons due to electric and magnetic fields, the electric field strength required for charge neutrality is given by

\[ \mathbf{E} = -\nabla \phi - \frac{\partial \mathbf{A}}{\partial t} = -\frac{\nabla p_e}{n_e e} - \mathbf{u} \times \mathbf{B} \]  

(2.57)

where \( \phi \) is the electrostatic potential, \( \mathbf{A} \) is the changing magnetic vector potential, \( n_e \) is the electron density, and \( p_e \) is the electron pressure given by \( p_e = n_e k_B T_e \). Substituting this into Faraday’s Law then gives

\[ \frac{\partial \mathbf{B}}{\partial t} = -\nabla \times (\mathbf{u} \times \mathbf{B} + \frac{\nabla p_e}{n_e e}) \]  

(2.58)

Thus a magnetic field is generated when \( \nabla \times (\nabla p_e / n_e) \neq 0 \) or \( \nabla n_e \times \nabla T_e \neq 0 \). In response to the magnetic field, many electrons will move across the field, thus some of the absorbed energy is lost. Forslund and Brackbill [41] found that this loss is about 30%. Dahmani and Kerdja [38] compared a flux limited model with a model that includes the effect of magnetic fields and found that a magnetic field of 76 T agrees with the flux limit range \( 0.05 \leq f \leq 0.08 \), which was found upon comparison with data.

2.3 Hydrodynamics

The hydrodynamics describing the evolution of laser-heated material can be described by solving the differential form of the fluid equations enforcing the conservation of mass,
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momentum, and energy (in one-dimension):

\[
\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x}(\rho u) = 0
\]  
(2.59)

\[
\frac{\partial}{\partial t}(\rho u) + \frac{\partial}{\partial x}(P + \rho u^2) = 0
\]  
(2.60)

\[
\frac{\partial}{\partial t}\left(\frac{3}{2}\rho c_s^2 + \frac{1}{2}\rho u^2\right) + \frac{\partial}{\partial x}\left[\rho u\left(\frac{u^2}{2} + \frac{5}{2}c_s^2\right) + Q\right] + I_A \delta(x - x_0) = 0
\]  
(2.61)

where \(\rho\) is the mass density:

\[
\rho = n_i A_m + n_e m_e \simeq n_e A_m p / Z
\]  
(2.62)

\(u\) is the plasma velocity, and \(P\) is the plasma pressure which is determined from the equation of state for the target material, \(I_A \delta(x - x_0)\) is a delta function describing the absorption of laser energy at position \(x = x_0\), the absorption surface, and the speed of sound in a plasma \(c_s\) is found by

\[
c_s = \left(\frac{Z + 1}{A m_p k_B T_e}\right)^{1/2}
\]  
(2.63)

These equations describe the system in the plasma corona, ablation zone, and compressed solid with the appropriate boundary conditions enforced. Continuity is required at the corona-ablation zone, and ablation front-compressed target interfaces.

To obtain a simple analytic solution, the coronal plasma is assumed isothermal. The solutions to the time-dependent Equations (2.59-2.61) then give an isothermal rarefaction wave propagating back into the dense plasma at the ion-acoustic speed and ions accelerating into the vacuum. Ablation pressure and mass ablation rate are then determined using the appropriate boundary conditions for the ablation zone. Here energy transport occurs mainly by electron thermal conduction. The solution for shock propagation can be described using the Rankine-Hugoniot relations [42] for mass, momentum, and energy conservation:

\[
\frac{V_0}{V} = \frac{D}{D - u}
\]  
(2.64)
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\[ P - P_0 = \frac{Du}{V_0} \]  \hspace{1cm} (2.65)

\[ E - E_0 = \frac{1}{2} P(V - V_0) \]  \hspace{1cm} (2.66)

where \( V \) is the specific volume, \( D \) is the shock speed, \( u \) is the change in particle speed across the shock front, and \( E \) is the internal energy. The subscript \( ^0 \) indicates the initial state of the shock and no subscript indicates the state behind the shock. This gives a set of three equations with five variables. The equation of state, \( P = P(V, E) \) for the material provides a fourth equation. The system can now be reduced to one equation with one free parameter which provides the locus of thermodynamic states (or Hugoniot curve) for a shock wave.
Chapter 3

Numerical Simulations

3.1 The Model

To quantify the effect of the value chosen for the flux limiter in numerical simulations, the hydrocode LTC (Laser Target Code) was used. This computer code was developed at the University of British Columbia by Peter Celliers [43] and is loosely based on the laser fusion code MEDUSA. LTC calculates the hydrodynamics of a laser heated target in one dimension by solving the plane parallel fluid equations for a single fluid where electrons and ions are assumed to have the same velocity to preserve charge neutrality. There is the option to perform one- or two-temperature calculations corresponding to electrons and ions having the same or different temperatures.

The target is divided into a number of cells at a particular initial temperature and density, and LTC follows the time evolution of each cell using the Lagrangian formalism, where time \( t \) and the Lagrangian coordinate \( m \) are independent variables. The definition of \( m \) (in kg/m\(^2\)) in terms of the density profile \( \rho(r, t) \) is given by [44]

\[
m(r, t) = \int_{R_i(t)}^{r} \rho(r', t) dr'
\]

where \( R_i(t) \) is the position of the free (rear) target surface at time \( t \), and \( r \) is the cell position with \( r \) and \( t \) in the laboratory frame.

The changes in the state of the system as it is heated and compressed are determined by following the laws of mass, momentum, and energy conservation. The physical
Chapter 3. Numerical Simulations

processes involved can be described by the set of fluid equations:

\[
\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x}(\rho u) = 0
\]  
(3.2)

\[
\frac{\partial}{\partial t}(\rho u) + \frac{\partial}{\partial x}(\rho u^2) = -\frac{\partial}{\partial x}(P_e + P_i)
\]  
(3.3)

\[
\frac{\partial}{\partial t}(\rho E_e) = -\frac{\partial}{\partial x} \left[ \rho u \left( E_e + \frac{P_e}{\rho} \right) \right] + \frac{\partial}{\partial x} \left( \kappa \frac{\partial T_e}{\partial x} \right) + \frac{\partial \epsilon_L}{\partial t} - g(T_e - T_i) \frac{\rho}{\rho_0} + u \frac{\partial P_e}{\partial x}
\]  
(3.4)

\[
\frac{\partial}{\partial t}(\rho E_i + \frac{\rho u^2}{2}) = -\frac{\partial}{\partial x} \left[ \rho u \left( E_i + \frac{P_i}{\rho} \right) \right] - \frac{\partial}{\partial x} \left[ \rho u \left( \frac{u^2}{2} \right) \right] + g(T_e - T_i) \frac{\rho}{\rho_0} - u \frac{\partial P_i}{\partial x}
\]  
(3.5)

LTC directly solves these equations for the mass density \(\rho\), fluid velocity \(u\), and material internal energies \(E_e\) and \(E_i\) (for electrons and ions respectively). \(\rho_0\) is the target solid density, \(\kappa\) is the thermal conductivity, \(g\) is the electron-ion coupling constant, and \(\epsilon_L\) the absorbed laser energy. Pressure has a component for electrons \((P_e)\) and ions \((P_i)\), and radiation pressure is negligible.

Fluid continuity and momentum conservation are represented in Equations (3.2) and (3.3). Equations (3.4) and (3.5) are used for energy conservation, one each for the electron and ion contributions. The energy equations are derived from the first law of thermodynamics. In Equation (3.4), the second term on the right represents heating due to thermal transport, the next term is the laser energy absorption rate (Section 2.1) followed by heat lost to the ions.

In numerical simulations, \(\kappa\) may be interpolated from data tables (such as the SESAME [45] or quotidian equation of state (QEOS) data libraries [46]) or using Spitzer’s equations for electrons [31] (from Equation (2.40)) and ions [4]:

\[
\kappa_e = 1.895 \times 10^{-10} \frac{(Z + 0.24)}{(1 + 0.24Z)} \frac{T_e^{5/2}}{Z \ln \Lambda} \quad \text{(W/Km)}
\]  
(3.6)

\[
\kappa_i = 4.3 \times 10^{-12} \frac{T_i^{5/2}}{Z \ln \Lambda}
\]  
(3.7)
LTC uses the harmonic mean of $\kappa_e$ to incorporate the flux limiter:

$$\kappa'_e = (\kappa_e^{-1} + \kappa_{elim}^{-1})^{-1}$$  \hfill (3.8)

where $\kappa_{elim}$ is the flux-limited electron conductivity. $\kappa'_e$ is then substituted for $\kappa$ in Equation (3.4).

For weakly coupled plasmas, $g = c_i/\tau_{eq}$ where $c_i$ is the ion specific heat and $\tau_{eq}$ is the equilibrium time. Spitzer [47] provides an expression for $\tau_{eq}$ for the case of weakly coupled, nondegenerate plasmas:

$$\tau_{eq} = \frac{3}{8\sqrt{2\pi}} \frac{m_i k_B T_e}{Z^3 n_i Z^3 e^4 \ln \Lambda} \left( \frac{k_B T_e}{m_e} \right)^{1/2}$$  \hfill (3.9)

For degenerate plasmas, Brysk [48] gives:

$$\tau_{eq} = \frac{3\pi m_i \hbar^3}{8 m_e^2 Z^3 e^4 \ln \Lambda} \left[ 1 + e^{-\mu/k_B T_e} \right]$$  \hfill (3.10)

where $\mu$ is the chemical potential and $\hbar = h/2\pi$ with $h$ as Planck’s constant. However, simple analytic expressions for $\tau_{eq}$ and $g$ are not available for strongly coupled plasmas. The assumption that electrons and ions have the same temperature $T$ is made for the simulations performed for this thesis. The fluid can then be described using one internal energy $E$ with pressure $P$, and the $g$ term is zero. A sample calculation is performed using the two-temperature model to show that this assumption is valid.

Once these principle variables are determined, other quantities related to the system may be obtained, in particular, the equation of state (EOS) which is expressed as

$$T_e = T_e(V,E_e), \quad P_e = P_e(V,E_e)$$  \hfill (3.11)

$$T_i = T_i(V,E_i), \quad P_i = P_i(V,E_i)$$  \hfill (3.12)

The ion number density is determined using Equation (2.9), from which the electron number density is found using Equation (2.8) where the average ionization state of a
cell is determined by interpolation of data from a table (such as SESAME or QEOS) for
the material used or by calculation with an atomic model. Other atomic and material
parameters such as conductivities and opacities are also found in the data table.

3.2 Results

3.2.1 Input Parameters

Numerical simulations were performed using both absorption methods discussed, i.e. the
traditional IB (referred to as the IB model in the remainder of this thesis) and EMS
models, for the range of possible flux limiters, $f$, between 0.03 and 0.6 to determine the
sensitivity of the shock characteristics to $f$, particularly the shock speed, predicted by
each model. If there is sufficient resolution, the value of $f$ for the system, within reason­
able uncertainty, may be estimated once the shock speed is known. It is also important
to determine if the IB results of the simulations performed are dependent on the value
of the fractional laser energy dump, $a$, at $n_{cr}$. This fraction is a designated percentage
of the laser energy reaching the cell which is then absorbed by that cell. Calculations
were performed for various values of $f$ with $a=1\%$, 10\%, 20\%, and 50\%. Maximum
laser irradiances of $2\times10^{13}$, $2\times10^{14}$, and $1\times10^{15}$ W/cm$^2$ were also compared at normal
incidence. The sensitivity to laser intensity is of interest because the absorbed energy
increases with laser power; therefore, the heat conduction changes and the dependence
on the flux limiter may change. Also, there is uncertainty associated with the laser power
measurement. The laser pulse was a p-polarized, 0.532 $\mu$m trapezoidal pulse with 100
ps rise and fall time, and 400 ps flat top. The corresponding laser pulse at an irradiance
of $2\times10^{14}$ W/cm$^2$ will be referred to as the standard pulse. This laser pulse is similar to
one of the pulse forms of the NOVA laser at Lawrence Livermore Laboratories.

An aluminum foil target was used for the simulations, which was thick enough that
the laser light did not burn through before 700 ps (this varied from 12-50 \( \mu \text{m} \)). The foil was divided into 3 zones (Figure 3.1). The first region was initially divided into very fine cells of uniform thickness (1364 cells were used, each \( \sim 2.4 \) \( \AA \)). The laser beam was incident on this zone and these cells expand quickly into the vacuum, increasing greatly in size. For these reasons it was important to have good resolution here in the initial stages of the plasma. The next region was geometrically zoned (the ratio 0.95 was used for 100 cells) and provides a smooth transition from the first zone to the third zone which was divided into much thicker cells of uniform thickness (0.04 \( \mu \text{m} \) was chosen, and the number of cells used was determined by the thickness of the foil). Ideally fine zoning would be used throughout, but resolution must be balanced with limitations on computational time. The cell sizes were chosen such that their size has no artificial effect on the results. The QEOS data tables [49] were used for material equation of state, conductivities, and ionization state. Electrons and ions were assumed to be at the same temperature. The target was initially at room temperature (300 K), and solid density (2700 kg/m\(^3\)).

3.2.2 Dependence of Shock Speed on Flux Limiter

The output of LTC provides the absorbed laser power, ionization, temperature, pressure, density, speed, and position of each cell in the target at specified time intervals. Figures 3.2 and 3.3 are examples of the graphical output for the standard pulse for the IB and EMS models, respectively. The laser beam is incident from the right side and the shock wave propagates to the left, into the target material. The corona is the first region seen from the right, and extends to the cell at or just below critical density, where laser energy deposition ends. In this case, IB absorption is weak through the majority of the expanding plasma, and a peak in absorbed laser intensity \( (I_A) \) is observed at \( n_{cr} \). With significant IB absorption, no peak is observed and \( I_A \) drops off before \( n_{cr} \). The
Figure 3.1: Schematic of a 3 zone scheme for a target. Zones 1 and 3 are uniformly divided, and zone 2 is geometrically scaled.
Figure 3.2: Sample of LTC output of the IB calculation for the standard laser pulse with 1% energy dump at $n_{tr}$, and $f=0.03$. rho is the cell density, and XI is the absorbed laser power. This snapshot is taken at 100 ps.
Figure 3.3: Sample of LTC output of the EMS calculation for the standard laser pulse with $f=0.03$. This snapshot is taken at 100 ps.
absorbed laser power curve for the EMS model shows the oscillations of the electric field. The ablation zone is the next region and the boundary between this region and the conduction zone occurs where the cell velocity goes from positive (to the right, out of the target) to negative (to the left, into the target). The absolute cell velocity is plotted here, so the boundary occurs at the velocity cusp. The shock front is seen as the sharp change in density, pressure, and temperature. From the output file the position of the shock front is estimated (Figures 3.4 and 3.5) and thus shock speed (Figures 3.6 and 3.7) can be determined. Examples of the temperatures and pressures reached at the shock front are given in Figures 3.8-3.11.

Notice that the shock does not reach a steady state in most cases. In the IB calculations it was observed that when the heat conduction is near classical the velocity reaches a maximum and then decreases even though the laser intensity is constant and at the maximum irradiance. As the material is ablated it expands rapidly into the vacuum; therefore, the distance between the cells where absorption occurs and the ablation front increases (Figures 3.12 and 3.13) and it takes longer for thermal conduction to transport the absorbed energy to the front and less energy reaches the target. In the IB results with no or little flux inhibition, this distance begins to increase at a slightly faster rate around 100 ps compared with higher inhibition (Figure 3.12) which results in a decrease in ablation pressure (Figure 3.14) and shock velocity. The sharp increase in this distance towards the end of the pulse is due to high energy attenuation by IB absorption before the critical density layer is reached; therefore, the cell which absorbs the greatest energy is before the critical density layer, unlike at earlier times where a peak in laser power absorption is seen at $n_{cr}$ where resonance absorption occurs. For high flux inhibition, this peak is also observed at $n_{cr}$. The critical density (from Equation (2.5)) is $3.89 \times 10^{21} \text{ cm}^{-3}$ for a 0.532 $\mu \text{m}$ laser. No sharp jump in the ablation zone thickness is observed for the EMS model (Figure 3.13) due to the difference in calculating absorption, and the
Figure 3.4: Shock position versus time for the standard trapezoidal pulse using the IB absorption model with various flux limiters and 1% laser energy dump at $n_{cr}$. 
Figure 3.5: Shock position versus time for the standard trapezoidal pulse using the EMS absorption model with various flux limiters.
Figure 3.6: Shock speed versus time for the standard trapezoidal pulse using the IB absorption model with various flux limiters and 1% laser energy dump at $n_{cr}$. 

[Graph showing shock speed versus time with various flux limiters and 1% laser energy dump at $n_{cr}$]
Figure 3.7: Shock speed versus time for the standard trapezoidal pulse using the EMS absorption model with various flux limiters.
Figure 3.8: Electron and ion temperature at shock front versus time for the standard trapezoidal pulse using the IB absorption model with various flux limiters and 1% laser energy dump at $n_{cr}$. 
Figure 3.9: Electron and ion temperature at shock front versus time for the standard trapezoidal pulse using the EMS absorption model with various flux limiters.
Figure 3.10: Shock pressure versus time for the standard trapezoidal pulse using the IB absorption model with various flux limiters and 1% laser energy dump at $n_{cr}$. 
Figure 3.11: Shock pressure versus time for the standard trapezoidal pulse using the EMS absorption model with various flux limiters.
Figure 3.12: Ablation zone thickness versus time for the standard trapezoidal pulse using the IB absorption model with various flux limiters and 1% laser energy dump at $n_{cr}$. 
Figure 3.13: Ablation zone thickness versus time for the standard trapezoidal pulse using the EMS absorption model with various flux limiters.
Figure 3.14: Pressure at ablation front versus time for the standard trapezoidal pulse using the IB absorption model with various flux limiters and 1% laser energy dump at \( n_{cr} \).
Figure 3.15: Pressure at ablation front versus time for the standard trapezoidal pulse using the EMS absorption model with various flux limiters.
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pressure at the ablation front smoothly increases until the laser intensity decreases. The noise seen in the velocity curves is due to the resolution of the cells. The actual position of the shock front is located within a cell which has a finite thickness.

Figure 3.16 shows that there is a significant difference, approximately 28% for $f=0.03$ and $\alpha = 1\%$, between the IB and EMS absorption methods. Figures 3.17-3.19 help to explain the cause of this discrepancy. Figures 3.17 and 3.18 illustrate that more laser energy is absorbed throughout the duration of the pulse (600 ps) for the IB model compared with the EMS model (about 68% for $f=0.03$ at 600 ps). This energy is absorbed mainly into cells with higher density (Figure 3.19). At early times absorption occurs at higher densities for the EMS model, but this is at low laser intensities when the pulse is still rising. Overall there is greater energy absorption into higher density material for the IB model. This causes greater mass ablation and pressure at the ablation front to conserve momentum (Figures 3.14 and 3.15), resulting in higher shock velocities (Figures 3.6 and 3.7).

Shock speed is greatly dependent on the value of flux limiter chosen (Figure 3.16) for $0.03 \lesssim f \lesssim 0.08$ for the IB model, and $0.03 \lesssim f \lesssim 0.15$ for the EMS model. For larger values of $f$, the shock speed is essentially constant; therefore, $f$ cannot be determined without significant uncertainty. For these $f$'s, heat conduction obeys the classical model. The energy absorption, cell density where this absorption occurs, and ablation pressure for different values of $f$ are compared in Figures 3.17, 3.19, and 3.14 for the IB model, and Figures 3.18, 3.19, and 3.15 for the EMS model. As the flux inhibition is decreased, the total laser power absorbed reaches a limiting value; therefore, the dependence on $f$ decreases.
Figure 3.16: Shock speed at 600 ps versus flux limiter for the standard pulse. A comparison between the IB and EMS absorption models is shown, as well as between various fractions of laser energy deposition at $n_{cr}$. The experimental shock speed is also provided.
Figure 3.17: Energy absorption coefficient versus time for the standard trapezoidal pulse using the IB absorption model with various flux limiters and 1% laser energy dump at $n_{cr}$. The laser irradiance is normalized to 1, where $I_{max}=2\times10^{14} \text{ W/cm}^2$. 
Figure 3.18: Energy absorption coefficient versus time for the standard trapezoidal pulse using the EMS absorption model with various flux limiters. The laser irradiance is normalized to 1, where $I_{\text{max}} = 2 \times 10^{14}$ W/cm$^2$. 
Figure 3.19: Density of cell with maximum laser power deposition versus time using the standard pulse for $f=0.03, 0.1$ for both models, with 1% laser energy dump at $n_{cr}$. 

Chapter 3. Numerical Simulations
3.2.3 Effect of Fractional Dump on IB Results

The IB method requires the use of a fractional laser energy deposition at $n_{cr}$ which is set by the user. The actual amount of absorption at this surface is not known, and would change in time, so it is important to determine how sensitive the numerical results are to this additional free parameter. This fraction represents resonance absorption and is expected to be important in regimes where IB absorption is not very strong. When IB absorption dominates, most of the laser energy is absorbed before the critical density layer so it does not matter what fraction of the remaining energy (if any) is allowed to be absorbed at $n_{cr}$. Figures 3.20-3.22 illustrate the variation in shock speed that can be obtained by changing the amount of energy absorbed at $n_{cr}$, and the results are summarized in Figure 3.16 which shows that in the regime where shock speed is sensitive to the flux limiter ($f \sim 0.08$), the speed is also sensitive to the value of the energy dump, particularly for $f=0.03$ where the shock speed for 1% and 10% dump differ by as much as 22%. No dependence is seen once free-streaming flow is predicted. The greater shock speed with higher fractional dump is not surprising since more energy is being absorbed into the system.

There is a much larger difference in the energy absorbed between the results for different energy dump using $f=0.03$ (Figure 3.23) than using $f=0.6$ (Figure 3.24). For example, the difference is $\sim 135\%$ between 1% and 50% dump for $f=0.03$, and only $\sim 1\%$ for $f=0.6$. Most of the laser energy is absorbed, and by the end of the pulse this absorption occurs before the critical density with larger $f$. This indicates that IB absorption is stronger with less flux inhibition, and the fractional energy dump at $n_{cr}$ is not as important. Also, the density at the absorption surface (Figures 3.25 and 3.26) is very similar for high values of $f$, but quite different for small values. This may explain why the sensitivity of shock speed with $f$ decreases with an increase in the absorbed
Figure 3.20: Shock speed versus time for the standard trapezoidal pulse using the IB absorption model with $f=0.03$ for various fractional laser energy dumps, $n$, at $n_{cr}$. 
Figure 3.21: Shock speed versus time for the standard trapezoidal pulse using the IB absorption model with $f=0.06$ for various fractional laser energy dumps, $n$, at $n_{cr}$. 
Figure 3.22: Shock speed versus time for the standard trapezoidal pulse using the IB absorption model with $f=0.1$ for various fractional laser energy dumps, $n$, at $n_{cr}$.
energy at the $n_{cr}$ layer.

Figures 3.27 and 3.28 show that the temperature at the absorption surface is essentially independent of $f$ for greater energy dump and independent on the energy dump for low flux inhibition. The IB absorption coefficient, $\alpha$, depends on electron temperature, number density and plasma ionization state. For large values of $f$, $T_e$ does not vary greatly, the density is fairly constant with $f$, and the plasma is almost fully ionized for much of the laser pulse; therefore, $\alpha$ is essentially the same for large $f$ which leads to the leveling off of the $v-f$ curve. The fluctuations seen results from the resolution of the cells chosen.

These results also illustrate that there is no unique set of parameters for a particular shock speed using the IB method. For example, a shock speed of $3 \times 10^6$ cm/s was measured by Bob Cauble using the NOVA laser facility at Lawrence Livermore Laboratory. This speed is predicted using the IB model with 1% dump and $f \approx 0.045$, and with 10% dump and $f \approx 0.03$. The EMS model gives this result with $f \approx 0.09$. This is a key result for those who use simulations with the IB absorption method. The IB absorption model should be used with caution as the physics may not be correctly represented. The EMS model has the advantage that there is only one free parameter set by the user, $f$.

3.2.4 Effect of Laser Beam Characteristics on Flux Inhibition

Laser Irradiance

The purpose of this investigation was to determine what laser intensity regime would provide the best resolution for determining the flux limiter. Calculations were performed for both absorption models, choosing 1% dump at $n_{cr}$ for the IB model. The results (Figures 3.29 and 3.30) show that a larger range of shock speeds is covered, and there is an improvement in resolution for higher laser irradiances. Consider the IB case.
Figure 3.23: Energy absorption coefficient versus time for the standard pulse using the IB model with $f=0.03$ and various fractional energy dump.
Figure 3.24: Energy absorption coefficient versus time for the standard pulse using the IB model with $f=0.6$ and various fractional energy dump.
Figure 3.25: Density at absorption surface versus time for the standard pulse using the IB model with $f=0.03$ and various fractional energy dump.
Figure 3.26: Density at absorption surface versus time for the standard pulse using the IB model with $f=0.6$ and various fractional energy dump.
Figure 3.27: Temperature at cell with maximum laser power absorption versus time for the standard pulse using IB absorption with 1% dump at $n_{cr}$. 
Figure 3.28: Temperature at cell with maximum laser power absorption versus time for the standard pulse using IB absorption with 10% dump at $n_{cr}$. 
Figure 3.29: Shock speed at 600 ps versus flux limiter for the IB absorption model for different laser irradiances, using 1% dump at $n_{cr}$.
Figure 3.30: Shock speed at 600 ps versus flux limiter for the EMS absorption model for different laser irradiances.
For $I_{max} = 2 \times 10^{13}$ W/cm$^2$ the shock speed is constant after $f \approx 0.05$ with a range of $1.45 \times 10^6 \lesssim v \lesssim 1.74 \times 10^6$ cm/s, while for $I_{max} = 1 \times 10^{15}$ W/cm$^2$ the speed reaches the free-streaming value at $f \approx 0.15$ with a range in speeds of $3.67 \times 10^6 \lesssim v \lesssim 7.31 \times 10^6$ cm/s. Similarly for the EMS case, shock speed is independent on $f$ for values $\lesssim 0.05$ with a range of $1.22 \times 10^6 \lesssim v \lesssim 1.64 \times 10^6$ cm/s for $I_{max} = 2 \times 10^{13}$ W/cm$^2$, but for $I_{max} = 1 \times 10^{15}$ W/cm$^2$ this occurs for $f \approx 0.15$ giving speeds of $2.38 \times 10^6 \lesssim v \lesssim 5.4 \times 10^6$ cm/s.

The aim of this work is to compare experimental data with the simulations to determine the best flux limiter. Uncertainty is always associated with measurements, so it is important to know how sensitive the numerical results are to laser intensity. A set of calculations was performed for the three intensity regimes examined using both the IB (Figures 3.31-3.33) and EMS (Figures 3.34-3.36) models and assuming a 10% uncertainty in intensity. The uncertainty in intensity for a specific value of $f$ results in similar uncertainty in speed variation for the different intensity regimes. For example, using EMS with $f = 0.04$, the shock speed varies by 4% for both $I_{max} = 2 \times 10^{13}$ W/cm$^2$ and $1 \times 10^{15}$ W/cm$^2$. For the standard pulse with $v = 3 \times 10^8$ cm/s, $f \approx 0.045 \pm 0.002$ using the IB model (Figure 3.32). When the EMS method is implemented (Figure 3.35), the uncertainty in $f$ is much higher, $f \approx 0.090_{-0.003}^{+0.010}$. If there is 5% uncertainty in the shock speed measurement as well, the uncertainty in the IB and EMS estimates of $f$ increases to $f \approx 0.045_{-0.007}^{+0.006}$, and $f \approx 0.090_{-0.022}^{+0.095}$, respectively.

### 3.2.5 Pulse Rise Time

The standard pulse used in the simulations has a 100 ps rise time. However, the pulse at Lawrence Livermore Laboratory has a rise time of about 150 ps. A comparison, using the IB model, between the two pulses was made to see how much of an effect a longer rise time with the same maximum intensity has on the shock characteristics. The shock front produced by the longer pulse was slower accelerating during the rise of the pulse,
Figure 3.31: Shock speed at 600 ps versus flux limiter for the IB absorption model showing the effect of a 10% variation in laser irradiance of $2 \times 10^{13}$ W/cm$^2$, using 1% dump at $n_{cr}$. 
Figure 3.32: Shock speed at 600 ps versus flux limiter for the IB absorption model showing the effect of a 10% variation in laser irradiance of $2 \times 10^{14}$ W/cm$^2$, using 1% dump at $n_{cr}$. 
Figure 3.33: Shock speed at 600 ps versus flux limiter for the IB absorption model showing the effect of a 10% variation in laser irradiance of $1 \times 10^{15}$ W/cm², using 1% dump at $n_{cr}$. 
Figure 3.34: Shock speed at 600 ps versus flux limiter for the EMS absorption model showing the effect of a 10% variation in laser irradiance of $2 \times 10^{13}$ W/cm$^2$. 
Figure 3.35: Shock speed at 600 ps versus flux limiter for the EMS absorption model showing the effect of a 10% variation in laser irradiance of $2 \times 10^{14}$ W/cm$^2$. 
Figure 3.36: Shock speed at 600 ps versus flux limiter for the EMS absorption model showing the effect of a 10% variation in laser irradiance of $1 \times 10^{15}$ W/cm$^2$. 
but once the laser light reached the maximum constant intensity the shock speed was the same (Figure 3.37).

### 3.2.6 Pulse Shape

For completeness, a comparison was made between the standard trapezoidal pulse and a triangular pulse with 100 ps rise and fall times, $I_{\text{max}}=2 \times 10^{14}$ W/cm$^2$ using the EMS model. The $x-t$ and $v-t$ graphs for the triangular pulse are given in Figures 3.38 and 3.39. No significant improvement in sensitivity to $f$ was observed (Figure 3.40).

### 3.2.7 Two-Temperature Calculation

The previous simulations were performed using a one-temperature code, which assumes electrons and ions are at the same temperature. To show that this is a valid assumption, a sample two-temperature calculation was performed, which treats electrons and ions separately. LTC uses a free parameter called the electron-ion coupling constant $g$ [50] to describe the equilibrium between electrons and ions. For silicon, Ng et al. [50] provide experimental evidence that $g \approx 10^{16}$ W/m$^3$K. The shock position and speed in aluminum using values of $g=10^{16}$, $10^{17}$, and $10^{18}$ W/m$^3$K are provided in Figures 3.41,3.42 using the EMS model with $f=0.03$ and a laser peak irradiance of $2 \times 10^{14}$ W/cm$^2$. They show that at the shock front, electrons and ions are essentially in equilibrium since the difference between the shock speed at 600 ps predicted by the one- and two-temperature models differs by $\sim$3% at most. Thus the one-temperature approximation is reasonably valid.
Figure 3.37: Shock speed versus time for the IB absorption model showing the comparison between 100 ps and 150 ps rise times for the trapezoidal laser pulse.
Figure 3.38: Shock position versus time for the EMS absorption model with various flux limiters, and using the triangular laser pulse.
Figure 3.39: Shock speed versus time for the EMS absorption model with various flux limiters, and using the triangular laser pulse.
Figure 3.40: Shock speed at 600 ps versus flux limiter for the EMS absorption model showing comparison between the standard trapezoid pulse and the triangular pulse.
Figure 3.41: Shock position versus time for the standard trapezoidal pulse using the EMS absorption model with $f=0.03$ providing a comparison of the one- (1-T) and two-temperature (2-T) models.
Figure 3.42: Shock speed versus time for the standard trapezoidal pulse using the EMS absorption model with $f=0.03$ providing a comparison of the one- ($1-T$) and two-temperature ($2-T$) models.
Chapter 4

Conclusions

A measurement of shock speed in laser-generated plasmas can be used to estimate the best value for the flux limiter $f$ to use in numerical simulations as the speed is highly sensitive to $f$ for $f \lesssim 0.08$ at $2 \times 10^{13}$ W/cm$^2$, and for $f \lesssim 0.15$ at $1 \times 10^{15}$ W/cm$^2$. Predictions of shock speed are independent of $f$ for larger values and the heat flux is almost classical. The EMS model of laser energy absorption is the preferred method to use because it only requires the use of one free parameter, $f$. Using this method, a shock speed of $3 \times 10^6 \pm 5\%$ cm/s generated by a trapezoidal laser pulse with irradiance $\sim 2 \times 10^{14} \pm 10\%$ W/cm$^2$ can be predicted using $f \approx 0.090^{+0.095}_{-0.022}$. The traditional IB model does not have a unique solution for shock speed. This speed can be predicted using 1\% energy dump at $n_{cr}$ with $f \approx 0.045^{+0.006}_{-0.007}$ or using 10\% dump with $f \approx 0.03$.

It was also shown that the effect of flux inhibition is greater for higher irradiances. However, use of a triangular pulse does not provide any advantage in determining $f$. Also, 50\% increase in pulse length does not change the results.

It is important to realize that the form of flux inhibited flow is not derived rigorously from physical processes, but is a phenomenological method used in an attempt to match theoretical calculations with experimental data. This leads to the risk that when the flux limit is used in simulations, and even in agreement with data, the physics is not accurately described, and other processes that are occurring may not be included, resulting in the wrong conclusions. For example, several investigations [9, 51] using thin foils resulted in a double peaked ion current distribution observed in charge collectors, such as Faraday
cups, rather than the usual single peak. This charge separation was only produced in simulations when strong flux inhibition \( f \leq 0.1 \) was used. The first ion peak is due to electrostatic acceleration of ions in the blow off region and the second peak is due to the expansion of the remaining plasma. It was believed that reduced heat flux resulted in an anomalously large fraction of energy being converted to kinetic energy of the fast ions, and less energy being conducted to the second front that propagates into the target. Gitomer and Henderson [52] later showed that the two-peaked velocity distribution can be modeled without using inhibited heat flow. They modified the LASNEX code to include the ponderomotive force and modified the hot electron temperature, finding that the shape of the curve depended on the hot electrons. Results suggested that fast ions were due to a ponderomotive force which stops the flow from one side of the foil. This behavior is similar to the effect of flux inhibition. This work provides an alternative explanation for only one phenomenon that was offered as evidence of reduced heat flow, and does not discount other effects which may truly result in reduction of heat flow.

4.1 Future Work

Radiation transport [53] was neglected because of time limitations and it is expected to have little effect except perhaps at high irradiances. Calculations should be performed to verify this using the code HYRAD [54] developed at the University of British Columbia. This code is coupled with LTC to include the effects of radiation emission, absorption, and transport, and calculates the new ionization of the target material.

Three parameters result in radiation emission in laser-heated plasmas: free-free collisions (bremsstrahlung emission), free-bound collisions (recombination emission), and relaxation or bound-bound transitions (line emission). Bremsstrahlung emission is the opposite process as inverse bremsstrahlung absorption, resulting from the acceleration of
an electron in the electrostatic field of an ion. Recombination emission occurs when a
free electron is captured by an ion into a bound energy level and a photon is emitted.
These two processes give rise to continuum radiation. Line emission results from the
transition of an electron between energy levels of an ion and a photon is released.

The next step is to collect accurate experimental data for the shock speed produced
by a laser with similar characteristics as that used in these simulations. Shock speed can
be determined by using a step target with a step of thickness $\Delta x$. The shock breakout
time at the two edges is measured. The speed $s$ then $\Delta x/\Delta t$ where $\Delta t$ is the difference is
breakout times. Once the speed is known, the appropriate flux limiter can be determined
from the $v-f$ curves.

Finally, once the value of the flux limiter is known, this result can be compared with a
nonlocal electron thermal transport model. The advantage of the nonlocal method is that
it does not depend on the free parameter $f$. Bell et al. [14] show that heat flow at any
position in the plasma is not simply determined by the local plasma state, the assumption
of the flux-limited Spitzer-Härm model, but is a function of the velocity distribution over
a few electron mean free paths. The nonlocal model solves the Fokker-Planck equation
for the electron distribution function. Heat flux is calculated using a convolution of the
Spitzer-Härm heat flux $Q_{SH}$ with a delocalization function:

$$Q(x) = \int dx'Q_{SH}(x')W(x, x')$$  \hspace{1cm} (4.1)

The delocalization kernel proposed by Luciani et al. [55] has the form

$$W(x, x') = [2\lambda(x')]^{-1}\exp[-|\int_{-x}^{x} dx''n(x'')/\lambda(x')n(x')|]$$  \hspace{1cm} (4.2)

where $\lambda(x')$ is the effective range of electrons at temperature $T_e(x')$:

$$\lambda(x') = a(Z + 1)^{1/2}\lambda_o(x')$$  \hspace{1cm} (4.3)
\( \lambda_0(x') \) is the electron mean free path, and the constant \( a \) is adjusted to fit results from Fokker-Planck simulations. Luciani \textit{et al.} found \( a \approx 32 \). This comparison will lead to improvements in the models and should further understanding of the physics of electron thermal transport.


