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Vol. I. Theory of the Coronal Plasma in Laser-Fusion Targets

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PREFACE

This report was originally prepared as a series of lectures for the Summer School of Theoretical Physics, Les Houches, France, in July of 1980 (to be published in the proceedings of the Summer School). I would like to thank Professor R. Balian, the Director of the Les Houches Summer School, and Drs. J.-C. Adam and R. Pellat, the organizers of the July 1980 session, for their hospitality and excellent organizational abilities.

I also want to thank my colleagues at Livermore and at the École Polytechnique for the many long hours they spent discussing with me the subjects covered in this report. I hope that my friends and associates found these exchanges as stimulating as I did. I especially would like to thank J. Fabre at the École Polytechnique for being my host in France, and J. Arons for helping to clarify some of the puzzles that arose while writing these lectures.

I am aware that the present report has tended to emphasize work done by my immediate colleagues and by myself, perhaps at the expense of equally good research done elsewhere. For this I apologize; it does not reflect on the quality of the work I have neglected to discuss, but rather on my own finite ability to assimilate the large literature on laser-plasma interactions.

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PHYSICS OF LASER FUSION
Vol. I. Theory of the Coronal Plasma
in Laser Fusion Targets

ABSTRACT

This monograph deals with the physics of the coronal region in laser fusion targets. The corona consists of hot plasma which has been evaporated from the initially solid target during laser heating. It is in the corona that the laser light is absorbed by the target, and the resulting thermal energy is conducted toward cold high-density regions, where ablation occurs. The topics to be discussed are theoretical mechanisms for laser light absorption and reflection, hot-electron production, and the physics of heat conduction in laser-produced plasmas. An accompanying monograph by H. Ahlstrom (Vol. II) reviews the facilities, diagnostics, and data from recent laser fusion experiments.

1. PHYSICAL CONDITIONS IN THE CORONA

When a laser beam shines on an initially solid target, a plasma (ionized gas) is rapidly formed on the front surface of the target. Thereafter, the propagation of the laser light is modified by the free electrons in the plasma. The dispersion relation for the laser light is

$$\frac{k_L^2 c^2}{\omega_L^2} = 1 - \frac{\omega_p^2}{\omega_L^2} \quad (1)$$

where k_L and ω_L are the wavenumber and frequency of the laser, and $\omega_p = (4\pi e^2 n_e / m_e)^{1/2}$ is the "plasma frequency" and is proportional to the square root of the electron density n_e . Thus, as the laser propagates into the plasma from lower toward higher electron densities, its wave vector becomes smaller and smaller. Finally, at the "critical density" $n_c \equiv m_e \omega_L^2 / 4\pi e^2$, the laser light can no longer propagate; for $n_e > n_c$ the wavenumber becomes imaginary and the light wave is evanescent.

Thus the critical density

$$n_c = 1.1 \times 10^{21} \left(\frac{1 \mu\text{m}}{\lambda_L} \right)^2 \text{ cm}^{-3} \quad (2)$$

is the classical reflection point for a laser of wavelength λ_L in a plasma. Since the laser cannot penetrate to densities above n_c , all laser light absorption and reflection processes occur at electron densities $n_e \leq n_c$ in the "underdense plasma." In typical laser-fusion targets, the laser light is absorbed at densities either at or close to the critical density, by processes that are discussed in Section 2. The light may also be prematurely scattered or distorted in the underdense plasma, as discussed in Section 3. The thermal energy deposited in the plasma by the laser light absorption is then conducted toward colder, higher-density regions by electron thermal conduction, which will be discussed in Sections 4 and 5. Figure 1 summarizes these effects. Section 6 analyzes the dependence of laser-plasma interactions on laser wavelength.

We begin by giving some order-of-magnitude estimates of the prevailing physical conditions in laser-produced plasmas. These simple physical estimates can be justified in more detail by a variety of analytic and numerical models. They are useful because the temperature, density, and velocity profiles of the background plasma play a crucial role in determining which mechanisms of absorption and heat transport will prevail in a given experimental situation.

1.1 TYPICAL TEMPERATURES NEAR THE CRITICAL SURFACE

Assume that most of the absorption of laser energy occurs at a density either at or close to n_c . We show in Section 2 that this is frequently the case. Then in a steady-state situation the absorbed energy must be carried away by heat conduction, which in a plasma is dominated by electrons. For the purpose of this simple discussion we estimate the electron heat flux by the phenomenological expression

$$q_e = -\ell(T_e, m_e)^{1/2} (n_e T_e) \text{ erg/cm}^2 \text{ s}^{-1} \quad (3)$$

where T_e is the electron temperature in energy units, and $\ell < 1$ is a "flux limit." Physically, Eq. (3) means that the electrons can transport their thermal energy density $n_e T_e$ at an effective speed ℓ ($\ell = 1$ m).

The balance between the rate of absorption of laser energy and the rate of heat conduction near n_c then becomes

$$L_{\text{abs}} = \ell(T_e, m_e)^{1/2} (n_e T_e) \quad (4)$$

FIG. 1. Schematic picture of electron density profile in a laser-produced plasma, showing location of major plasma-physics coupling processes. The "critical density" $n_c = m_e \omega_L^2 / 4\pi e^2$ is the highest electron density accessible to laser light of frequency ω_L .

where I_{abs} is the absorbed laser intensity. This expression can be solved for the electron temperature T_e

$$T_e \text{ (keV)} = 0.6 \left[\frac{1}{1} \left(\frac{I_{abs}}{10^{14} \text{ W/cm}^2} \right) \left(\frac{\lambda_L}{1 \mu\text{m}} \right)^2 \right]^{2/3} \quad (5)$$

For example if $I_{abs} = 10^{14} \text{ W/cm}^2$ and $f = 0.03$ (as described in Section 4), Eq. (5) gives electron temperatures of 6.4 keV for $\lambda_L = 1.06 \mu\text{m}$, and 0.9 keV for $\lambda_L = 1.4 \mu\text{m}$.

Detailed discussions of the usefulness of Eq. (3) for the electron heat flux may be found in Refs. [1]-[3]. For the purposes of the present order-of-magnitude estimates it is sufficient to point out that Eq. (5) is not valid if a separate distribution of "hot" or "suprathermal" electrons can carry much more heat flux than the thermal electrons, or if there are very large discontinuities in the plasma density or velocity at the critical-density surface. These issues are addressed in Section 2.2.

1.2 TYPICAL VELOCITIES IN THE CORONA

In laser-heated initially solid targets, the inward flow of thermal energy toward cooler, high-density regions produces a compensating outward flow of plasma kinetic energy toward lower-density regions. The resulting outward velocity of the plasma toward the laser plays an important role in determining the mix of absorption, reflection, and transport processes that can occur.

A simple estimate is that the blow-off velocity is roughly equal to the local sound speed, for $n_e \gg n_i$

$$v \approx c_s = \lambda \sqrt{Z T_e / M_i} = 3 \times 10^7 \left[\left(\frac{Z}{A} \right)^{1/2} \left(\frac{T_e}{1 \text{ keV}} \right)^{1/2} \right] \text{ cm/s} \quad (6)$$

where (Z/e) and $M_i = A m_p$ are the ionic charge and mass, respectively. For typical conditions this gives blow-off velocities between 10^7 and 10^8 cm/s in the corona. At very high laser intensities where many suprathermal

electrons are produced, the suprathermal electrons can at times determine the temperature to be used in Eq. (6), as described in Ref. [4].

Several different types of analytic theories [1,5,6] indicate that the blow-off velocity should be at least sonic [Eq. (6)] at the critical density surface, and that it should then rise to supersonic values $v > c_s$ for densities $n_e < n_c$. While this behavior is confirmed in computer-hydrodynamics calculations where the laser intensity is steady over relatively long periods of time [1], hydrodynamic simulations of shorter-pulse experiments under some conditions show regions of underdense plasma where the flow is slightly subsonic [7]. Thus Eq. (6) should not be regarded as more than a crude estimate.

1.3 GRADIENT SCALE LENGTHS IN THE CORONA

Because the coronal electron temperature is typically a few keV or larger, thermal conduction in the underdense plasma tends to maintain a nearly uniform electron temperature distribution in space. Thus typical electron temperature scale lengths $L_T \equiv T_e / |dT_e/dx|$ are very long, for $n_e < n_c$. The ion temperature in the corona is generally much smaller than the electron temperature: ions cool due to expansion but have no heat sources in the corona (ion thermal conduction is small, and electron-ion coupling is usually negligible). Thus the ion temperature and pressure are usually small relative to the electrons, and gradients in ion temperature have little physical impact.

By contrast, gradient scale lengths for density and velocity are usually determined by the hydrodynamics of the blow-off. For short laser pulses the scale lengths are roughly $c_s \tau_L$, where τ_L is the laser pulse length and c_s the coronal sound speed [5,6]. On the other hand, if the laser pulse is so long that $c_s \tau_L$ becomes larger than the effective radius of curvature due to irradiation or target geometry, the geometrical divergence of the flow will determine the density and velocity scale lengths [1,8]. These two extremes are illustrated in Fig. 2. In terms of an effective radius of curvature R_{eff} , the density and velocity gradient lengths are roughly

$$L_i \approx \min(c_s \tau_L, R_{\text{eff}}) \quad (7)$$

$$R_{\text{eff}} = \begin{cases} R_{\text{sphere}}, & \text{spherical geometry} \\ 2R_{\text{focal spot}}, & \text{planar geometry} \end{cases}$$

That is, for $c_s \tau_L > R_{\text{eff}}$ the background geometrical divergence due to target or focal-spot geometry dominates the outflow structure.

These arguments suggest a criterion for the maintenance of planar geometry in a given laser-target experiment. Even if the initial solid target is planar, the blow-off will remain roughly planar only if $c_s \tau_L < 2R_{\text{focal spot}}$, or

$$R_{\text{focal spot}} > 150 \mu\text{m} \left[\left(\frac{Z}{A} \right)^{1/2} \left(\frac{T_e}{1 \text{ keV}} \right)^{1/2} \left(\frac{\tau_L}{1 \text{ ns}} \right) \right] \quad (8)$$

Thus experiments using long pulses require large laser spots, if plane geometry is to be maintained. This constraint can be an important one, since—for fixed laser intensity—Eq. (8) implies that the required laser energy increases as the cube of the pulse length τ_L .

Geometrical considerations such as these have led McCrory et al. [9] to numerically model planar target experiments by using sectors of spheres with radius $R_{\text{eff}} = 2R_{\text{focal spot}}$.

FIG. 2a. Isodensity contours typical of short-pulse experiments. The density scale length L_n is determined by the laser pulse length for both planar and spherical targets: $L_n \approx c_s \tau_L$.

FIG. 2b. Isodensity contours typical of long-pulse experiments. The density scale length is limited by the geometrical divergence of the plasma flow. For planar targets the effective radius of curvature R_{eff} is roughly twice the radius of the focal spot.

1.4 PLASMA PARAMETERS

In the corona of a laser fusion target, the electron Debye length $\lambda_{De} \equiv (T_e / 4\pi n_e e^2)^{1/2}$ is typically much smaller than the laser wavelength. Numerically,

$$\lambda_{De} = 7.4 \times 10^{-7} \text{ cm} \left[\left(\frac{T_e}{1 \text{ keV}} \right) \left(\frac{10^{21} \text{ cm}^{-3}}{n_e} \right) \right]^{1/2} \quad (9)$$

Thus for a CO_2 laser, a Nd-glass laser, and a KrF laser with coronas at typical temperatures T_e , we have at the critical density:

TABLE 1. Typical values of the Debye length λ_{De} at three different laser wavelengths λ_L .

λ_L	T_e	λ_{De}
10.6 μm	20 keV	0.3 μm
1.06 μm	10 keV	2.0 μm
0.76 μm	0.5 keV	1.3 μm

From a theoretical point of view, much of standard plasma kinetic theory assumes that there are many particles in a Debye sphere. That is, $N_D \gg 1$ where

$$N_D = \frac{4\pi}{3} n_e \lambda_{De}^3 = 1.7 \times 10^3 \left[\left(\frac{T_e}{1 \text{ keV}} \right)^{3/2} \left(\frac{10^{21} \text{ cm}^{-3}}{n_e} \right)^{1/2} \right] \quad (10)$$

This condition is most likely to become marginal for the higher densities and lower temperatures typically produced by short-wavelength lasers. For example, if $\lambda_L = 0.26 \mu\text{m}$, the temperature at $n_e = 5n_c$ might be $< 50 \text{ eV}$, in which case $N_D < 2$.

Another assumption of standard kinetic theory [10] is that typical gradient scale lengths l of temperature, density, and velocity are large compared to the electron-ion mean free path

$$\lambda_{mfp} \approx 30 \mu\text{m} \left[\left(\frac{T_e}{1 \text{ keV}} \right)^2 \left(\frac{10^{21} \text{ cm}^{-3}}{Z n_e} \right) \left(\frac{10}{n \Lambda} \right) \right] \quad (11)$$

Here $n\Lambda$ is the usual Coulomb logarithm ($n\Lambda = 5-10$) and Z is the ionic charge. The discussion in the preceding section shows that the condition $l \gg \lambda_{mfp}$ is rather easy to violate in the underdense plasma, particularly for short-pulse experiments, where $l \approx c \tau_L$ can be quite small.

Recently, large dc magnetic fields have been observed in laser-produced plasmas. (This subject is discussed in detail in Section 5.) A measure of the effect of a magnetic field on electron heat transport is given by the ratio of the electron-ion mean free path to the electron gyroradius, λ_{mfp}/ρ_1 , where $\rho_1 = (4\pi n_e)^{1/2} c/eB$. If $\rho_1 \ll \lambda_{mfp}$, then the effective step-size describing diffusive electron heat flow becomes ρ_1 instead of λ_{mfp} , in the direction perpendicular to the magnetic field B . Magnetic fields as large as a megagauss have been observed experimentally. Under these conditions the ratio λ_{mfp}/ρ_1 is

$$\frac{\lambda_{mfp}}{\rho_1} \approx 40 \left[\left(\frac{T_e}{1 \text{ keV}} \right)^{3/2} \left(\frac{10^{21} \text{ cm}^{-3}}{Z n_e} \right) \left(\frac{10}{n \Lambda} \right) \left(\frac{B}{10^6 \text{ gauss}} \right) \right] \quad (12)$$

Thus for low densities and Z s that are not too large, it is possible to have $\lambda_{mfp}/\rho_1 \gg 1$. Equivalently, the product of the electron cyclotron frequency and the electron-ion collision time is large: $\Omega_{ce} \tau_{ei} \gg 1$. The magnetic field can then cause a potentially large inhibition of the electron heat flow.

2. MECHANISMS FOR LASER LIGHT ABSORPTION

Essential to almost all laser fusion schemes is the concept that inward momentum is imparted to the imploding target by the ablation of a hot gas, rather than by the momentum of the laser photons themselves. The reason is clear: for a given laser energy E , the ratio of available momentum to available energy is c^{-1} for the laser photons. By contrast, this ratio is $2v^{-1}$ if the laser light is first absorbed by a plasma that then ablates at a velocity v . Thus for a hot gas with a (nonrelativistic) ablation velocity v , the momentum per unit energy is $c/v \gg 1$ times that available for photon pressure alone.

As a consequence, efficient absorption of laser light energy by the coronal plasma is crucial to laser fusion. In fact, the ultimate efficiency of a laser fusion power plant is directly proportional to the fraction of incident laser energy absorbed by the target. In the present section we discuss the three absorption mechanisms presently thought to be most important: 1) inverse Bremsstrahlung (sometimes called "classical" or "collisional" absorption), 2) resonance absorption, and 3) absorption due to ion-acoustic turbulence.

2.1 INVERSE BREMSSTRAHLUNG ABSORPTION

Inverse Bremsstrahlung absorption is due to electron-ion collisions in the plasma. If plasma electrons are subject to momentum-changing collisions as they oscillate back and forth in the laser electric field, the laser light wave feels an effective damping. Crudely speaking, the energy damping rate of the light wave ν_l can be obtained from the damping of the oscillatory kinetic energy of plasma electrons moving in the laser electric field E_l

$$\nu_l \left(\frac{E_l^2}{8\pi} \right) = \nu_{ei} \left[\frac{n_e}{2} m_e \left(\frac{eE_l}{m_e \omega_l} \right)^2 \right] .$$

Solving for ν_l and recalling that $\omega_p^2 \equiv 4\pi e^2 n_e / m_e$, we obtain

$$\nu_l = \left(\frac{\omega_p^2}{\omega_l^2} \right) \nu_{ei} = \left(\frac{n_e}{n_c} \right) \nu_{ei} . \quad (13)$$

The scaling of the electron-ion collision frequency ν_{ei} can be estimated as follows: The collision frequency is $\nu_{ei} \approx n_i v_{ie}$ for an electron of thermal velocity v_{ie} colliding with background ions of number density n_i . The collision cross section is $\sigma \approx \pi b^2$, where the distance of closest approach b is given by the balance of potential and kinetic energy:

$$\frac{Ze^2}{b} = \frac{1}{2} m_e v_{ie}^2 ,$$

for ions of charge Ze . Solving for b and substituting, we find

$$\nu_{ei} \approx n_i \left(\frac{2Ze^2}{m_e v_{ie}^2} \right)^2 v_{ie} \propto \frac{n_i Z^2}{T_e^{3/2}} \propto \frac{n_e Z}{T_e^{3/2}} . \quad (14)$$

In laser fusion applications, it is generally most relevant to know the spatial damping rate of an incident laser light wave of given frequency ω_l . If we include the effect of electron-ion collisions, the laser light-wave dispersion relation (1) becomes

$$\frac{k^2 c^2}{\omega_l^2} = 1 - \frac{\omega_p^2}{\omega_l(\omega_l + i\nu_{ei})} \approx 1 - \frac{\omega_p^2}{\omega_l^2} + \left(\frac{i\nu_{ei}}{\omega_l} \right) \left(\frac{\omega_p^2}{\omega_l^2} \right) \quad (15)$$

where we have expanded assuming $\nu_{ei}/\omega_l \ll 1$. This can be solved for k ,

$$k = \pm \frac{\omega_l}{c} \left(1 - \frac{\omega_p^2}{\omega_l^2} \right)^{1/2} \left\{ 1 + i \left(\frac{\nu_{ei}}{2\omega_l} \right) \left(\frac{\omega_p^2}{\omega_l^2} \right) \left[1 - \frac{1}{\omega_p^2/\omega_l^2} \right] \right\} .$$

The spatial damping rate of wave energy, κ_{ib} , is then twice the imaginary part of k :

$$\kappa_{ib} \approx \frac{\nu_{ei}}{c} \frac{\omega_p^2}{\omega_L^2} \frac{1}{(1 - \omega_p^2/\omega_L^2)^{1/2}} \quad (16a)$$

$$\kappa_{ib} \propto \frac{Z n_e^2}{T_e^{3/2} (1 - n_e/n_c)^{1/2}} \quad (16b)$$

Inverse Bremsstrahlung is thus strongest for low temperatures, high densities, and high-Z target materials.

A more rigorous calculation of the inverse Bremsstrahlung absorption coefficient can be made by using kinetic theory, to take into account the actual distribution function of the electrons and the actual positions of the ions. A clear review of this approach may be found in Ref. [11]. The net absorption is shown to depend on the correlation function of the background ions. Highly correlated ions created, for example, by ion-acoustic turbulence can result in enhanced absorption. This mechanism is discussed in Section 2.3. On the other hand, a thermal distribution of ions and a Maxwellian distribution of electrons produce the following inverse Bremsstrahlung absorption coefficient [12] for $\omega_L \gg \omega_{pi}$:

$$\kappa_{ib} = (2\pi)^{1/2} \left(\frac{16\pi}{3} \right) \frac{Z n_e^2 e^6 / n \Lambda}{c (m_e T_e)^{3/2} \omega_L^2 (1 - n_e/n_c)^{1/2}} \quad (17)$$

$$\Lambda \equiv \frac{\nu_{te}}{\omega_p \rho_{min}} \quad , \quad \rho_{min} \equiv \max \left[\frac{Z e^2}{T_e} \cdot \frac{\hbar}{(m_e T_e)^{1/2}} \right] \quad (18)$$

Thus a uniform plasma of length L and density $n_e < n_c$ would produce a one-way absorption fraction

$$\alpha_{abs} = 1 - \exp(-\kappa_{ib} L) \quad (19)$$

For weak absorption, $\kappa_{ib} L \ll 1$ and $\alpha_{abs} \approx \kappa_{ib} L$. For strong absorption, $\kappa_{ib} L \gg 1$ and $\alpha_{abs} \rightarrow 1$. So the absorption fraction is linear in $\kappa_{ib} L$ for small absorption, and then saturates when $\kappa_{ib} L$ is large. In general long plasma lengths L enhance the absorption.

Results for inverse Bremsstrahlung absorption, in an inhomogeneous plasma are more complicated, since in Eq. (17) n_e , T_e , and $n \Lambda$ may all depend on position. In reality, however, T_e varies much more slowly in space than n_e , and the variation of $n \Lambda$ is sufficiently slow that it can be neglected. Thus the dominant spatial dependence in Eq. (16a) or (17) comes from $n_e(x)$.

The classic reference on the subject of inverse Bremsstrahlung in a density gradient is Ginsburg [13]. Ginsburg shows by integrating Eq. (16a) from $n_c = 0$ to $n_e = n_c$, and then back to $n_e = 0$, that the two-way absorption fraction for a linear density profile with scale length L is

$$\alpha_{abs} = 1 - \exp \left(- \frac{32}{15} \frac{\nu_{ei}(n_c)}{c} L \right) \quad (20)$$

Thus in integrated form, all the density-dependent factors in Eq. (16a) combine to yield the numerical coefficient $32/15$ for a linear profile, with the collision frequency ν_{ei} now evaluated at the critical density n_c . The increased absorption for low T_e , high n_e , and long L is retained in a density gradient.

Because of the dependence of κ_{ib} on $(n_e/n_c)^2 (1 - n_e/n_c)^{-1/2}$ in Eqs. (16a) and (17), a significant fraction of the inverse Bremsstrahlung absorption comes from the region very near the critical density n_c . But as we discuss below, there are several important processes which can steepen the density profile near n_c and which may thus eliminate much of the region where inverse Bremsstrahlung would be expected to be

FIG. 3. Model density profile for calculation of inverse Bremsstrahlung absorption in a steepened profile. The density is assumed to rise infinitely steeply above a density βn_c , $\beta < 1$.

strongest. This point was emphasized by Kruer several years ago [14]. To illustrate the effect, consider an electron density profile which is exponential up to a density $n_e = \beta n_c$, $\beta \leq 1$, and which then abruptly rises to $n_e = n_c$ with infinite slope, as illustrated in Fig. 3. The one-way absorption coefficient is given by

$$\int \kappa_{ih} dx = \alpha \int_{n_e=0}^{n_e=\beta n_c} \frac{\left(\frac{n_e(x)}{n_c}\right)^2}{\left(1 - \frac{n_e(x)}{n_c}\right)^{1/2}} dx \quad (21a)$$

$$\alpha = \frac{\nu_{ei}(n_c)}{c} \approx 3 \times 10^{-21} n_c Z / n \lambda \left(\frac{1 \text{ keV}}{T_e}\right)^{3/2} \text{ cm}^{-1} \quad (21b)$$

with density profile $n_e(x) = \beta n_c \exp(-x/L)$. The integral (21a) is straightforward and yields

$$\int_{n_e=0}^{n_e=\beta n_c} \kappa_{ih} dx = \left(\frac{4}{3} \alpha L\right) F(\beta) \quad (22a)$$

$$F(\beta) = 1 - (1 + \beta/2)(1 - \beta)^{1/2} \quad (22b)$$

The function $F(\beta)$ is plotted in Fig. 4. It is equal to 1 when $\beta = 1$, but it then falls quickly toward zero for $\beta < 1$. For example when $\beta = 1/2$ (maximum density $n_e = n_c/2$), the absorption is only 12% of its value for $\beta = 1$. The same physical effect is seen for linear [14] or other types of density profiles: inverse Bremsstrahlung absorption can be strongly diminished by profile steepening near the critical density surface.

In addition to profile modification there are two other nonlinear effects which can decrease the inverse Bremsstrahlung absorption coefficient for high laser intensities. In the first effect the electron distribution function is assumed to remain Maxwellian as in Eq. (17), but the electron-ion Coulomb collision frequency is modified by the electrons' coherent oscillatory motion in the laser electric field. Thus rather than having a typical mean-square velocity v_{te}^2 , the electron velocity instead has an effective dispersion

$$(v^2)_{\text{eff}} = v_{te}^2 + \left(\frac{eE_L}{m_e \omega_L}\right)^2 \quad (23)$$

FIG. 4. The factor $F(\beta)$ from Eq. (22), describing the reduction of inverse Bremsstrahlung optical depth in a steepened density profile. The parameter β is defined in Fig. 3.

where E_L is the laser electric field. Because the absorption coefficient scales as $\kappa_{ib} \propto |E_L|^{-2} \propto [\omega_p(\omega_L)]^{-2}$, at low laser intensities the coefficient κ_{ib} in Eq. (17) is replaced by

$$\kappa_{ib} \rightarrow \frac{\kappa_{ib}}{1 + \frac{3}{2} \left(\frac{eE_L}{m_e \omega_L v_{te}} \right)^2} \quad (24)$$

for $(eE_L/m_e \omega_L) < v_{te}$. This effect was originally discussed by Silin [15], and has been reviewed in Refs. [16] and [17].

A second nonlinear modification to inverse Bremsstrahlung was discussed recently by Langdon [18] and is due to the additional fact that the electron distribution function can at times be far from Maxwellian when the laser intensity is high. The rate of energy gain of an electron from the laser light is proportional to $\nu_{ei}(eE_L/m_e \omega_L)^2$. This is to be compared with the rate the electrons can share energy with other electrons to form a Maxwellian, $\nu_{ee} \approx \nu_{ei}/Z$ (Here $\nu_{ee} \approx \nu_{ei}/Z$ is the electron-electron energy sharing rate [19]). Thus the electrons gain energy from the laser faster than their distribution can become Maxwellian when

$$Z \left(\frac{eE_L}{m_e \omega_L v_{te}} \right)^2 > 1 \quad (25)$$

From a kinetic-theory point of view, the inverse Bremsstrahlung absorption rate is proportional to the number of electrons in the distribution which have their electron-ion collision frequency $\nu_{ei}(x)$ just equal to the laser frequency ω_L . These are the electrons whose oscillatory motions in the wave field are most strongly

FIG. 5. Non-Maxwellian electron distribution function produced by inverse Bremsstrahlung absorption [18] when $Zv_o^2 v_{te}^2 > 1$. The two curves show self-consistent distributions for $Zv_o^2 v_{te}^2 \ll 1$ (the Maxwellian) and $Zv_o^2 v_{te}^2 = 6$.

perturbed by collisions. Since $v_e(v) \propto v^{-1}$, and generally $v_e(v_e) \ll \omega_p$, those electrons responsible for the absorption have relatively low velocities compared with v_{te} . However, when Eq. (25) is satisfied, there is a general lack of these low energy electrons relative to their population in a Maxwellian, as shown in Fig. 5. Thus the amount of collisional absorption decreases below the value derived for a Maxwellian in Eq. (17), when $Zv_o^2 v_{te}^2 > 1$.

The quantitative decrease in absorption due to Langdon's mechanism is given to within a few per cent by

$$\frac{\kappa_{\text{th, actual}}}{\kappa_{\text{th, linear}}} = \exp \left(-\frac{v_w^2}{2v_{te}^2} \right), \quad v_w = Zv_o^2 v_{te}^2, \quad (26)$$

where v_w is that electron velocity for which $v_e(v_w) = \omega_p$

$$\frac{1}{2} m_e v_w^2 = 8 e V \left[Z \left(\frac{n_e}{n_c} \right) \left(\frac{1 \mu\text{m}}{\lambda_l} \right) \left(\frac{r n V}{10} \right) \right]^{2/3} \quad (27)$$

For realistic laser parameters and moderate- Z targets, the inverse Bremsstrahlung absorption can be reduced by as much as a factor of 2 below the classical value [Eq. (17)] for a Maxwellian.

In summary, inverse Bremsstrahlung absorption is due to electron-ion collisions. It is strongest for low electron temperatures, high densities, high- Z targets, and long density scale lengths. It is sensitive to details of the density profile near critical, and can be diminished considerably if the profile falls steeply at the critical surface. Finally, at high laser intensities and/or for target $Z \gg 1$, nonlinear effects can decrease the inverse Bremsstrahlung absorption coefficient below its usual linear value.

2.2 RESONANCE ABSORPTION

In resonance absorption [13,20], laser light obliquely incident on a plasma density gradient can excite resonant plasma oscillations at the critical density surface, as shown in Fig. 6. When the driven electrostatic plasma oscillation is damped either by collisional or collisionless processes, energy is in effect transferred from the laser electric field to the thermal energy of the plasma. The resulting absorption persists even when the classical collision frequency ν_{ei} is small. Thus resonance absorption tends to dominate over inverse Bremsstrahlung for high plasma temperatures (high laser intensities), longer laser wavelengths (lower values of n_e), and short plasma scale-lengths (low values of $k_{\parallel}$).

Figure 6 illustrates this process in more detail. For light incident on a plasma density gradient at an angle θ to the normal, Snell's law and Eq. (1) can be combined to show that the classical turning-point for the light wave occurs at a density $n_e(\theta) = n_c \cos^2 \theta$. If the electric field of the incident laser light is in the plane formed by $\mathbf{k}$ and ∇n_e , then at the turning-point the local electric field points in the direction of ∇n_e . Some of this electric field can tunnel in to the critical-surface region, where it can resonantly drive electron plasma oscillations (Langmuir waves).

It is clear from Fig. 6 that if the laser electric field $\mathbf{E}$ is perpendicular to the plane of $\mathbf{k}$ and ∇n_e , then at the turning point there will be no component of $\mathbf{E}$ along ∇n_e , and hence no driven Langmuir waves. This polarization-dependence of resonance absorption can also be seen from the following simple argument. Maxwell's equations for waves in a plasma can be rewritten in the form

$$\nabla \cdot (\epsilon \mathbf{E}) = 0 \quad (28a)$$

$$\epsilon^2 + 1 - \frac{\omega_p^2}{\omega^2} = 1 - \frac{n_e}{n_c} \quad (28b)$$

On the other hand, the electron charge-density perturbation δn_e of the Langmuir wave obeys

$$-4\pi e \delta n_e = \nabla \cdot \mathbf{E} = -\left(\frac{\mathbf{E} \cdot \nabla \epsilon}{\epsilon}\right) \propto \left(\frac{\mathbf{E} \cdot \nabla n_e}{\epsilon}\right) \quad (29)$$

Thus if $\mathbf{E} \cdot \nabla n_e = 0$, $\delta n_e = 0$ and there are no driven Langmuir waves. This geometry is frequently called "s" polarization. On the other hand, if $\mathbf{E} \cdot \nabla n_e \neq 0$ ("p" polarization), δn_e can be large when $\epsilon \rightarrow 0$ (i.e., when $n_e \rightarrow n_c$), and a large Langmuir wave is driven in the critical-surface region.

FIG. 6. Schematic configuration for resonance absorption of p-polarized light, where $\mathbf{E}$ is in the plane of $\mathbf{k}_{\parallel}$ and ∇n_e .

The actual amount of resonance absorption in p polarization depends on the angle of incidence θ . The electric field must tunnel the distance between $n_c = n_e \cos^2 \theta$ (the turning point) and $n_e = n_e$. If θ is too large, the wave must tunnel too far and the Langmuir wave is not driven efficiently. On the other hand, if θ is too small, the component of E parallel to ∇n_e at the turning point will be small (or vanishes at $\theta = 0$), and once again the plasma wave amplitude will be small. Thus the optimum angle for resonance absorption is intermediate. Quantitatively, the absorption can be parametrized by the function $\tau \equiv (\omega_L/c)^{1/3} \sin \theta$ in a linear density profile of scale length L . The peak in absorption at intermediate angles of incidence, for p polarization, is then clearly seen from the function $\phi(\tau)$ shown [13] in Fig. 7. The absorption fraction is equal to $(1/2)\phi^2(\tau)$, so Fig. 7 shows that absorption maximizes at about 60% for $\sin \theta \approx 0.8(\omega_L/c)^{1/3}$.

The actual amount of absorption at the peak of $\phi(\tau)$ is independent of the strength of damping felt by the resonant Langmuir wave. As with many resonant systems, when damping is weak the width of the resonance peak for the electric field of the Langmuir wave, $E(x)$, is very small, but its height is high. By contrast, the width of $E(x)$ for strong damping is large but the maximum height is small. Thus the integral under $E^2(x)$ does not depend on the damping rate, as long as nonlinear phenomena such as wavebreaking [21] or extreme profile steepening [22] are not dominant.

If a strong dc magnetic field is also present in the underdense plasma, resonance absorption can be enhanced for some angles of incidence and some polarizations relative to B . These effects are described in Rets. [23] [26].

FIG. 7. The function $\phi(\tau)$ characterizing resonance absorption [13]. The total absorption at an angle of incidence θ is equal to $(1/2)\phi^2(\tau)$, where $\tau \equiv (L\omega_L/c)^{1/3} \sin \theta$.

Particle-in-cell computer simulations of resonance absorption at high laser intensities ($I_L \lambda_L^2 > 3 \times 10^{14} \text{ W}\mu\text{m}^2 \text{ cm}^2$) tend to show collisionless heating of electrons traveling through the resonant plasma oscillation from high densities toward low densities [22,27]. These electrons see a nearly stationary electric field, and can then be strongly accelerated down the gradient. The heated electron distribution is approximately Maxwellian, for reasons that have been addressed only recently [28]. The average hot-electron energy T_h of this hot Maxwellian has been estimated as

$$\frac{T_h}{1 \text{ keV}} = 13 \left(\frac{T_c}{1 \text{ keV}} \right)^{1/4} \left[\left(\frac{I_L}{10^{15} \text{ W/cm}^2} \right) \left(\frac{\lambda_L}{1.06 \mu\text{m}} \right)^2 \right]^{0.39} \quad (30)$$

in Ref. [27], and as

$$\frac{T_h}{1 \text{ keV}} = 7 \left(\frac{T_c}{1 \text{ keV}} \right)^{1/3} \left[\left(\frac{I_L}{10^{15} \text{ W/cm}^2} \right) \left(\frac{\lambda_L}{1.06 \mu\text{m}} \right)^2 \right]^{1/3} \quad (31)$$

in Ref. [22]. Here T_c is the temperature of the cold background electrons, I_L is the laser intensity, and λ_L is the laser wavelength. (The reader is cautioned that I have multiplied the value of T_h given in Ref. [27] by 3/2, because Refs. [22] and [27] did not use the same definition of T_h in terms of the slope of the two-dimensional distribution function.)

We derived in Eq. (5) a different expression for the temperature as a function of $I_L \lambda_L^2$, which applies at low intensities where hot electrons do not dominate the energy flux. In that derivation the cold electron density was known, and the electron temperature was found as a function of the absorbed intensity. By contrast, in the high-intensity regime $I_L \lambda_L^2 > 10^{15} \text{ W}\mu\text{m}^2 \text{ cm}^2$, the hot-electron temperature is known from Eq. (30) or (31), and the absorbed intensity determines the density of hot electrons, n_h . In this case the absorbed laser energy is carried away dominantly by the hot electrons, so that instead of Eq. (4) we have

$$I_{\text{abs}} = \alpha_{\text{abs}} I_L = f n_h m_e (T_h/m_e)^{3/2}, \quad (32)$$

where α_{abs} is the absorption fraction and f the flux limit for the *hot* electrons. Thus the fraction of hot electrons at the critical density is

$$\frac{n_h}{n_c} = \frac{\alpha_{\text{abs}}}{2f} \left(\frac{v_o^2}{v_{te}^2} \right), \quad (33)$$

where we have defined $v_o^2 \equiv T_h/m_e$, and

$$v_o \equiv e E_L / m_e \omega_L, \quad (34)$$

is the oscillatory velocity of an electron in the laser field E_L .

We recall that v_{te} is given, for example, by Eq. (31), and that

$$\frac{v_o^2}{v_{te}^2} = 4 \left(\frac{I_L}{10^{16} \text{ W/cm}^2} \right) \left(\frac{\lambda_L}{1 \mu\text{m}} \right)^2 \left(\frac{1 \text{ keV}}{T_c} \right). \quad (35)$$

Then Eqs. (33)–(35) can be combined [22] with some algebra to find the fraction of hot electrons at the critical density

$$\frac{n_h}{n_c} = \frac{0.4 \alpha_{\text{abs}}}{f} \left(\frac{v_o}{v_{te}} \right). \quad (36)$$

For example for an absorption fraction of 30% and an f for suprathermal electrons equal to about 0.6, we find that n_h is 20% of critical when $v_o^2/v_{te}^2 \approx 1$. The condition $n_h \approx n_c$ is particularly easy to satisfy for high-power lasers of wavelength longer than one micron, which are therefore expected to produce coronas dominated by hot electrons.

In addition to hot-electron generation, computer simulations have shown the importance of profile-steepening in resonance absorption [22,29]. Figure 8, from Ref. [29], illustrates this phenomenon. Laser light obliquely incident from the left on an initially linear density profile ($\omega_1 t = 0$) drives a resonant Langmuir wave at the critical surface. The electric field pressure $E^2/8\pi$ of the plasma oscillation has pushed plasma out of the way by the time $\omega_1 t = 168$, causing a locally steep profile at n_c and a local density hump below n_c . By $\omega_1 t = 420$ the density hump has propagated away toward lower densities, leaving a profile whose density rises by more than a factor of 2 in a distance of less than one laser wavelength.

Such profile steepening has important consequences. First, it makes resonance absorption less sensitive to the angle of incidence of the laser, because there is now less distance to tunnel to reach the critical density surface. Second, it produces less energetic hot electrons than a non-steepened profile would, because the hot electron temperature $T_h \propto eEL$, and steepening reduces both the spatial extent L and the amplitude E of the resonant plasma wave. Third, profile steepening means that there is only a very small region near n_c where the parametric decay and oscillating two-stream instabilities [30] may occur; it thus strongly limits the effect of these two processes on absorption and preheat when the laser intensity is high.

Profile steepening at the critical surface is a complex process. From a theoretical point of view, at least three separate mechanisms control the amount of steepening expected. 1) Electric field pressure or "ponderomotive force": $E^2/8\pi$ of the laser light or of the plasma oscillation roughly balances $n_c U_c$ of the overdense plasma. Estimates of this effect [22] suggest that the density profile may be steepened to as high as several hundred times n_c for high-power CO₂ lasers ($\lambda_1 = 10.6 \mu\text{m}$). The dominant force is frequently due to the resonant plasma wave, which accelerates hot electrons down the density gradient during resonance absorption. To balance the momentum they carry away, the profile recoils near critical and steepens there. Since electrons have a higher ratio of momentum to energy than photons do, this mechanism can be more important than the ponderomotive force of the laser light itself, particularly at moderate laser intensities [31]. 2) Inhibited heat transport of the thermal (cold) electrons toward high densities can produce a density step at critical [1,32]. 3) If the density and temperature of hot electrons are high enough, a rarefaction shock [4] can form at n_c . In this situation the underdense plasma consists entirely of hot electrons.

As is clear from the above discussion, self-consistent profile steepening is a complex process. There is already a large theoretical literature on steepening mechanisms at low and high laser intensities. There are also a large number of experiments where steepening has been observed. I shall not attempt to discuss the details of profile steepening further in the present report.

2.3 ABSORPTION DUE TO ION-ACOUSTIC TURBULENCE

It was pointed out many years ago [33] that absorption of light by a plasma can be enhanced when coherent short-wavelength ion fluctuations are present. More recently [34,35], consequences of this process for high-power laser-plasma interactions have been explored. As we mentioned in Section 2.1, the usual kinetic-theory derivation of inverse Bremsstrahlung absorption [11,33] uses a thermal distribution of electrons and randomly distributed (uncorrelated) ions. However in the presence of coherent ion turbulence with fluctuation spectrum $\delta n_i(\mathbf{k})/n_i$, the energy damping rate of laser light of frequency ω_1 becomes

$$(\nu_E)_{\text{eff}} = \frac{\omega_1}{2} \sum_{\mathbf{k}} \left(\frac{\delta n_i(\mathbf{k})}{n_c} \right)^2 \text{Im} \left[\frac{1}{\epsilon(\mathbf{k}, \omega_1)} \right] (\hat{\mathbf{k}} \cdot \hat{\mathbf{E}}_L)^2 \quad (37)$$

where $\hat{\mathbf{E}}_L$ is the direction of the laser electric field, and $\text{Im}[1/\epsilon(\mathbf{k}, \omega_1)]$ is plotted [34] in Fig. 9. This figure shows that for short-wavelength ion fluctuations with $k \approx 1/(2\lambda_{De})$, the absorption rates between $0.4n_c$ and n_c can be substantial if the level of ion fluctuations $\delta n_i/n_i$ is large enough. The optimum density for absorption varies with the wavelength of the ion fluctuations.

Detailed consideration of this mechanism for laser-target experiments with pulse lengths $\tau_1 < 100$ ps (see Refs. [34] and [35]) indicates that for $n_e > n_c/3$ and $T_e > 5T_i$, ion turbulence might contribute up to an extra 20% to the total absorption. Absorption due to ion turbulence would add to the fraction absorbed by

FIG. 8. Development of a steep density profile during resonance absorption, from particle-in-cell computer simulations [29]. The angle of incidence is 24° , and the free-space incident laser intensity is 10^{16} W/cm 2 , $\lambda_L = 1.06$ μ m. The initial electron temperature is 4 keV.

FIG. 9. The imaginary part of $\epsilon^{-1}(k, \omega_1)$, as a function of plasma density [34]. The brackets denote an average over a uniform spectrum of ion fluctuations between $k = k_D/2$ and $k = k_D$.

resonance absorption, since the two processes maximize at different places in the plasma density profile. Reference [35] simulated absorption by ion-acoustic turbulence in long-pulse experiments ($\tau_L \approx 150$ ps, $\omega_L = 1.06 \mu\text{m}$), and found that ion-acoustic absorption can reach levels as high as 54% for a laser intensity of $7 \times 10^{14} \text{ W cm}^{-2}$. Electrons heated by short-wavelength ion-acoustic absorption are typically less hot than electrons heated by resonance absorption [Eqs. (30) and (31)]. The theoretical upper limit on the fluctuation amplitude in these circumstances is roughly $\delta n_i/n_i < 10\%$ for $k\lambda_{De} \approx 1-2$, and is due mainly to ion trapping.

The actual source and magnitude of short-wavelength ion turbulence in laser-produced plasmas remains somewhat hypothetical. Manheimer and his colleagues [34,35] have postulated that the turbulence is produced by heat-flow driven ion-acoustic instabilities. Faehl and Krueer [34] and Forslund and colleagues [36] suggested that ion-ion streaming instabilities could create the required 10% density-fluctuation levels. In both cases it is necessary to have a rather wide angular spectrum for the turbulence, because of the factor of $(\hat{k} \cdot \hat{E}_1)^2$ which appears in Eq. (37). It is my impression that there is not yet a good theoretical consensus on the expected magnitude and angular spectrum of ion turbulence. Fortunately for the theoreticians, ion turbulence has now been measured directly in several CO_2 laser-plasma experiments [32], and substantial enhancements in the ion fluctuation level over thermal noise have been seen. Although the angular spectrum has not yet been well-characterized experimentally, these measurements strongly increase the plausibility of absorption due to ion-acoustic turbulence.

If ion fluctuations are indeed present with the optimal characteristics $\delta n_i/n_i \sim 10\%$, $k\lambda_{De} \approx 1-2$, $T_e > 5T_i$, and a broad angular spectrum, they are in principle capable of increasing the expected absorption level considerably. The maximum absorption of about 20% due to ion turbulence calculated for short-pulse experiments [34] is quite comparable with the resonance absorption levels of 20-30% expected for mixed p and s polarizations. For long-pulse experiments, a more relevant comparison is the ratio of the optimal ion-acoustic absorption rate [Eq. (37)] to the inverse Bremsstrahlung absorption rate [Eq. (13)]:

$$\frac{(\nu_E)_{\text{eff}}}{(\nu_E)_{\text{IB}}} \cong \frac{10}{Z} \left(\frac{T_e}{1 \text{ keV}} \right)^{3/2} \left(\frac{\lambda_L}{1 \mu\text{m}} \right) \quad (38)$$

at $n_e = n_c/2$, $\delta n_i/n_i = 0.1$, $k\lambda_{De} = 1/2$, $T_e > 5T_i$, $(\hat{k} \cdot \hat{E}_L)^2 = 1/2$. This indicates that ion-acoustic absorption can be dominant for low-Z targets and laser wavelengths $\lambda_L \gtrsim 1 \mu\text{m}$, *provided* that a large enough region of plasma is present at the required optimum density.

In summary, in this section we have discussed three different laser light absorption mechanisms. Inverse Bremsstrahlung is due to classical electron-ion collisions. It increases with target Z, with plasma density, and with the length of underdense plasma, and it decreases at high plasma temperatures. It is thus strongest for low laser intensities, short laser wavelengths, and long laser pulses. Inverse Bremsstrahlung may be nonlinearly reduced at high laser intensity or high target Zs.

Resonance absorption, by contrast, tends to dominate at high laser intensities and short laser pulses. Because it depends on the excitation of resonant electron plasma oscillations at the critical surface, resonance absorption is both angle- and polarization-dependent. The theoretical absorption peaks at 50–60% in p polarization, implying that for mixed polarization the absorption should be about half this value. The critical density surface is thought to be steepened and rippled for high laser intensities $I_L \lambda_L^2 > 10^{18} \text{ W}\mu\text{m}^2/\text{cm}^2$. Under these conditions, resonance absorption creates a nearly-Maxwellian tail of hot electrons, and ejects them in the direction of decreasing plasma density.

Enhanced levels of ion-acoustic fluctuations can cause increased absorption. Under ideal conditions, ion turbulence might add up to 20% to the total absorption for short-pulse experiments. For long-pulse experiments, ion turbulence absorption can be larger than inverse Bremsstrahlung for low-Z targets and for laser wavelengths $\lambda_L > 1 \mu\text{m}$, *provided* that a long enough region of plasma exists near the optimum density. However there is not yet a theoretical consensus on the level and angular spectrum of ion turbulence expected in laser-produced plasmas.

3. STIMULATED SCATTERING AND FILAMENTATION: NONLINEAR EVOLUTION

In this section we discuss stimulated scattering, which might reduce laser light absorption below the values calculated in Section 2, and filamentation, which might enhance nonuniformities in the incident laser beam. In particular we shall describe: 1) stimulated Brillouin scatter, in which the laser light decays into an ion-acoustic wave and a backscattered light wave; 2) stimulated Raman scatter and two-plasmon decay, in which the laser light decays into a Langmuir wave and either a scattered light wave (Raman) or a second Langmuir wave (two-plasmon decay); and 3) filamentation or self-focusing, in which spatial inhomogeneities initially present on the incident laser beam can be amplified.

The linear instability theory of these processes is now quite well-developed. It is described in Ref. [38] and is reviewed extensively in Ref. [30]. However, theoretical understanding of the nonlinear evolution of stimulated scattering and filamentation is in a considerably earlier stage of development, and I shall attempt here to synthesize the major conclusions to date.

Because of their mathematical difficulty, nonlinear models are frequently very simplified; typical theoretical calculations include only one nonlinear effect at a time, in order to keep the problem soluble. Because of this, analytic nonlinear models are valuable when one nonlinear process is so strong as to dominate over other competing processes. On the other hand, computational studies can give insight into how several different nonlinearities compete with each other. Computer studies sometimes have the disadvantage of being very costly, however, particularly if one wants to calculate nonlinear evolution over the long length- and time-scales involved in large reactor-sized laser fusion targets.

The particular nonlinearities which dominate stimulated scattering and filamentation are sensitive to the size of the underdense plasma, L . Until recently, the majority of fusion-related experiments used laser pulse lengths $\tau_L < 100$ ps, laser spot sizes $< 100 \mu\text{m}$, and high laser intensities $I_L \lambda_L^2 > 10^{15} \text{ W cm}^{-2} \text{ cm}^2$. These parameters define the regime necessary for "exploding pusher" targets [39]. Under these conditions the length of the underdense plasma typically does not exceed 10–20 vacuum wavelengths ($L \lambda_L < 20$). With such a small underdense plasma the absorption physics is dominated by resonance absorption on a steepened critical surface, with perhaps some absorption by ion-acoustic turbulence. Inverse Bremsstrahlung, stimulated scattering, two-plasmon decay, and filamentation are relatively minor effects because they lack the space in which to grow.

By contrast, present and future experiments will emphasize more gentle, ablative implosions. The consequent longer laser pulses, larger focal spots, and lower laser intensities tend to create a larger region of underdense plasma with $L \lambda_L > 100$. Under these circumstances, many more physical processes can play a significant role. Inverse Bremsstrahlung, stimulated scattering, and filamentation may all become stronger, in a large underdense plasma whose size is many times the linear growth length for these instabilities.

3.1 STIMULATED BRILLOUIN SCATTERING

Stimulated Brillouin scattering (abbreviated as SBS) has the potential of causing large premature reflection of the incident laser light before it reaches the critical density, and of limiting the energy that can reach critical. The incident laser light decays into an ion-acoustic wave and a reflected light wave which can leave the plasma, $\omega_L \rightarrow \omega_{ia} + \omega_R$. The physical mechanism for this process is shown schematically in Fig. 10. An initial low-frequency density perturbation δn produces a transverse current $\propto \delta n E_L$, due to the oscillatory motion of electrons in the laser electric field E_L . This transverse current produces a reflected light wave with field E_R . The ponderomotive force ($E_L E_R / 8\pi$) of the incident and reflected light waves can in turn cause enhancement in the original density perturbation δn , and thus exponential growth, provided that the frequency and wavenumber of the three waves satisfy the conditions

$$\omega_L = \omega_R + \omega_{ia}, \quad k_L = k_R + k_{ia}.$$

FIG. 10. Feedback mechanism for growth of stimulated Brillouin scattering.

For the simplest case where

$$\omega_1 \gg \omega_{pe}, \quad k_1 \parallel k_R,$$

these matching conditions imply that

$$k_{ia} = 2k_1 \approx 2\omega_1/c,$$

$$k_R = -k_1 \approx -\omega_1/c,$$

$$\omega_R = \omega_1 - 2k_1 c_s.$$

The nonlinear evolution of SBS in short underdense plasmas was studied in Ref. [40]. Subsequent numerical simulations by Kruer, Valeo, and Estabrook [41] indicated that at high laser intensities where $2I_L/c \gg n_e T_e$, momentum deposition in the ions due to Brillouin-scattering causes the density profile to steepen. As a consequence of this profile steepening, the SBS reflectivity decreases after an initial transient. The process is illustrated in Fig. 11, from Ref. [41]. Such limitation of reflectivity in the short-pulse, high-intensity regime has been seen in Nd-glass laser experiments [42]. Recent high-power CO₂ laser experiments [43] are also in this regime.

If there is a very large region of underdense plasma, profile steepening occurs at a much slower relative rate and SBS can be much stronger, particularly when $2I_L/c \leq n_e T_e$. Under these conditions a wide range of physical mechanisms may determine the nonlinear evolution of SBS. Among these I shall discuss the following. 1) ion-wave saturation by trapping or wavebreaking, 2) nonlinear ion heating, 3) the effect of light reflected from the critical surface, 4) competition between SBS and inverse Bremsstrahlung absorption.

FIG. 11. Limitation of stimulated Brillouin scattering by profile steepening [41], when $2I_L/c \gg n_e T_e$. The laser comes from the left. The top figure shows the time evolution of the density profile. The bottom figure shows that as the profile steepens, the average transmission fraction increases towards unity due to the suppression of SBS. Parameters are $I_L = 5 \times 10^{15}$ W/cm², $\lambda_L = 1.06$ μ m, T_e ($t = 0$) = 1 keV, $M_i/m_e = 100$, initial $T_e/T_i = 5$.

3.1.1: Ion-Wave Saturation by Trapping or Wavebreaking

One way in which the growth of SBS might be saturated is to limit the size of the relative ion density perturbation, $\delta n_i/n_i$. Here n_i is the background ion density. It is clear physically that $\delta n_i/n_i$ must be < 1 . However in reality this limit may be considerably lower than unity, due to wavebreaking of the ion waves [44] or due to the fact that large-amplitude ion waves can cause ion trapping [45]. Wavebreaking and trapping are

most likely to occur at high laser intensities, and lead to estimates of $\delta n_i/n_i$ between 10 and 20%. When $\delta n_i/n_i$ is a (known) constant, the amount of SBS can be estimated from the following simple model [45]. Consider a plasma of length L and uniform density $n_e = Z n_i$, and assume that SBS is far above the convective threshold but is not an absolute instability (for example because the plasma outflow is supersonic in the lab frame, or because the ion wave damping is not very small). Then in one dimension the equations of evolution of the incident wave's vector potential A_i and the backscattered wave's vector potential A_b are

$$\frac{dA_i}{dx} = - \left(\alpha \frac{\delta n_i}{n_i} \right) A_b \quad , \quad (39a)$$

$$\frac{dA_b}{dx} = - \alpha \left(\frac{\delta n_i}{n_i} \right) A_i \quad , \quad (39b)$$

$$\alpha = \frac{\pi}{2} \frac{n_e}{n_c} \frac{1}{\lambda_L} \left(1 - \frac{n_e}{n_c} \right)^{-1/2} \quad . \quad (39c)$$

The sum and difference of Eqs. (39a) and (39b) may each be integrated from $x = 0$ to L , to obtain

$$\frac{A_i(L) + A_b(L)}{A_i(0) + A_b(0)} = \frac{A_i(0) - A_b(0)}{A_i(L) - A_b(L)} = \exp \left[- \alpha \left(\frac{\delta n_i}{n_i} \right) L \right] \quad . \quad (40)$$

For the case where $A_b(L) \ll A_i(L)$, this yields the SBS reflectivity [45]

$$R = \tanh^2 \left[\alpha \left(\frac{\delta n_i}{n_i} \right) L \right] \quad . \quad (41)$$

This expression shows that when the underdense plasma is large, the reflectivity can be substantial even when $\delta n_i/n_i$ is nonlinearly limited to a very small value. For example consider an underdense plasma with $n_e/n_c = 0.1$, and let us limit the ion wave amplitude to $\delta n_i/n_i = 0.01$, a value ten times smaller than the ion-trapping estimates mentioned above. Then Eq. (41) says the reflectivity R will be less than 50% only for $L/\lambda_L < 600$. This condition is relatively easy to violate with the large targets envisioned for fusion reactors, which have diameters up to a fraction of a centimeter, and which thus have underdense plasmas which can be thousands of wavelengths long for $\lambda_L = 1.06\text{-}\mu\text{m}$ laser light. We are hence motivated to look for other physical mechanisms which might nonlinearly limit SBS to a lower level of reflectivity.

3.1.2 Nonlinear Ion Heating

Computer simulations of strong SBS indicate [44] that a heated tail on the ion distribution is created in the nonlinear state. These nonlinearly heated ions are shown in Fig. 12. Physically, the energy source for such ion heating is the fraction $\omega_{ia}/\omega_L = 2v_{Te}/c$ of the incident light energy which is given to the ion wave during SBS.

Recently Krueer and colleagues [45,46] have suggested that nonlinear ion heating of this type might be able to limit SBS by increasing the damping rate of the ion waves. To see why increased damping could limit SBS, we solve again Eq. (39) for the reflectivity in a finite uniform plasma. However this time we are concerned with the case where the ion wave damping ν_i is large ($\nu_i/\omega_{ia} > 1$ kL_G, where L_G is the convective growth length for SBS). We assume that the ion wave amplitude $\delta n_i/n_i$ is driven by the ponderomotive force of the two light waves A_i and A_b , so that $\delta n_i/n_i$ is given by

$$\frac{\delta n_i}{n_i} = A_b A_i \left(\frac{Z k_{ia}^2 e^2}{2 m_e M_i c^2 \nu_i \omega_{ia}} \right) \quad . \quad (42)$$

FIG. 12. Nonlinear ion heating during SBS, from computer simulations by K. Estabrook. a) Ion phase space, showing large-amplitude ion acoustic waves. Parameters are $I_L = 2 \times 10^{16} \text{ W/cm}^2$, $\lambda_L = 1.06 \mu\text{m}$, $T_e = 2 \text{ keV}$, $T_e/T_i = 3$, 19.5° angle of incidence. b) Ion distribution function in the nonlinear state ($v_0/c = 0.1$, initial $T_e = 10 \text{ keV}$, initial $T_i = 3 \text{ keV}$, $n_e/n_c = 0.7$, $L/\lambda_L = 20$). Note the pronounced heated tail.

In general this corresponds to a regime of laser intensities lower than that characterizing ion wavebreaking (Section 3.1.1). The solution to Eqs. (39) and (42) was described by Tang [47]. The reflectivity R is obtained from the equation

$$R(t) - R = \frac{I_N}{I_L} \{ \exp[Q_{H1}(1 - R) - R] \} \quad (43)$$

$$Q_{H1} = \frac{\pi}{2} \left(\frac{L}{\lambda_L} \right) \left(\frac{n_e}{n_c} \right) \left(\frac{E_{L1}^2}{8\pi n_e T_e} \right) \left(\frac{\nu_i}{\omega_{ia}} \right)^{-1} (1 + 3T_i/ZT_e)^{-1} (1 - n_e/n_c)^{-1} \quad (44)$$

Here I_N is the intensity of the noise in the back-reflected wave at $x = 0$, and I_L is the incident laser intensity at $x = L$. The "quality factor" Q_{H1} is just the number of linear growth lengths for SBS in a homogeneous plasma of length L .

Figure 13 shows the SBS reflectivity R as a function of the "quality factor" Q_{H1} from Eq. (43). At small Q_{H1} , the reflectivity is at first low because the backscattered wave must be amplified from the noise level I_N to a point where it can start being comparable with I_L . This is the "linear" region ($R \ll 1$) of exponential amplification: $R \approx (I_N/I_L) \exp Q_{H1}$. The extent of this regime depends upon the noise level I_N . For larger values of Q_{H1} the amplified backscattered wave becomes comparable with the incident intensity I_L , and the rate of growth of R with Q_{H1} becomes slower than exponential due to pump depletion. At very large values of Q_{H1} the reflectivity saturates at $R = 1$, and the pump wave is entirely depleted.

To return to the subject of ion heating, we see from Eq. (44) that in this homogeneous plasma example, Q_{H1} decreases when the damping of the ion wave ν_i is large. Figure 13 then indicates that the SBS reflectivity should diminish as well. This is thought to be the case when ion-tail heating occurs, because the effective temperature of tail ions seen in computer simulations [44] is $T_e \approx M v_e^2$, exactly the temperature which produces strong ion Landau-damping of the ion-acoustic waves. The scenario proposed by Kruer [45,46] is thus that self-consistent ion tail heating due to SBS could reduce the SBS reflectivity in Eq. (43) by raising the ion wave damping rate ν_i to a value $\nu_i/\omega_{ia} \approx 1$. (An analogous mechanism for Raman scattering [48] will be discussed in the following section.)

It is instructive to ask for which laser and target parameters this SBS ion heating should be the most important. The ion tail heating rate, by the Manley-Rowe relations, is proportional to the rate of decrease of the incident light wave due to SBS: $(\omega_{L1} - \omega_L)(dI_L/dx)_{SBS}$. The heating rate is opposed by the conductive cooling rate due to free-streaming of the tail ions away from the heated region: $(d/dx)(n_e M v_e^2)$, where n_e is the

FIG. 13. SBS reflectivity R as a function of the "quality factor" Q_{H1} for a homogeneous plasma. Curves represent solutions of Eq. (43) for three different values of the noise level: $I_N/I_L = 10^{-4}$, 10^{-3} , and 10^{-2} .

density of heated ions in the tail. Thus a measure of the importance of ion tail heating is the ratio of the heating rate to the cooling rate

$$H = \frac{\text{heating rate}}{\text{cooling rate}} = \frac{\left(\frac{2k_1 c_s l_1}{\omega_1 - l_1} \right)}{\left(\frac{n_{th} M_i c_s^3}{l} \right)} = 2 \left(\frac{l}{l_1} \right) \left(\frac{E_1^2}{8\pi n_e T_e} \right) \left(\frac{n_e / Z}{n_{th}} \right) \quad (45)$$

Here l/l_1 is just Q_{H1} , the number of linear growth lengths l_1 in a homogeneous plasma of length l . Thus Eq. (44) for Q_{H1} implies that

$$H = \frac{1}{\lambda_L} \left(\frac{E_1^2}{8\pi n_e T_e} \right)^2 \left(\frac{n_e}{n_{th}} \right), \quad l_1^2 \lambda_L^2 l \quad (46)$$

Ion heating saturation of SBS should therefore be most important at high laser intensities and long laser wavelengths.

The above discussion of ion tail heating has concentrated on the case where the plasma is finite and homogeneous ($n = \text{const}$ and $v = \text{const}$). Braggian scattering in nonuniform plasmas differs in some important ways [49,50] from the homogeneous-plasma situation described above. In a plasma with a velocity gradient dv/dx , the frequency of ion-acoustic waves $\omega = k [c_s + v(x)]$ changes with position. The wave vectors of the three waves also depend on position, because the index of refraction falls as the electron density increases. Thus a given ion wave and backscattered light wave can only satisfy the frequency and wavenumber matching conditions within a small region around their resonant position x_R . If the ion wave damping $\gamma = \omega$ is very small, the local amplification rate [Q_{H1} in Eq. (44)] is large, but on the other hand the backscattered wave is only amplified within a short distance $\approx (v - \omega)/\omega$ of the resonant position x_R . [Here $l \equiv v/(dv/dx)$ for example.] In contrast if $\gamma = \omega$ is larger, the local amplification rate Q_{H1} drops but the growth distance $\approx (v - \omega)/\omega$ increases in the same proportion. This is illustrated in Fig. 14. The net result [49] is that the total linear amplification of the backscattered intensity in an inhomogeneous plasma is $\exp[(\pi/2)Q_{H1}(\omega - \gamma)]$ for the case of a

FIG. 14. Schematic amplification profiles for SBS in a strongly inhomogeneous plasma. The total amount of amplification, integrated over the whole plasma profile, is independent of the ion wave damping rate.

plasma of size L . This exponentiation is independent of the damping rate, because $Q_H \propto \omega_{ia}^{-1} \nu_i$. Thus although nonlinear ion heating increases the damping of the ion waves ν_i , Refs. [49] and [50] predict that the resulting higher damping does not change the linear growth of SBS in an inhomogeneous plasma. Reference [51] has corroborated this fact in computer simulations of strongly inhomogeneous underdense plasmas.

Thus we conclude that nonlinear ion heating will be effective at limiting the growth of SBS in situations when the plasma is effectively finite and homogeneous, so that Eqs. (43) and (44) apply. The condition for this to be the case is given by Pesme et al. [49] as

$$\frac{\nu_i}{\omega_{ia}} > \frac{\ell}{L}, \quad (47)$$

where ℓ is the geometrical length of the finite plasma and L is the local gradient length. If $\ell/L \ll 1$, finite-homogeneous SBS theory can apply even to weakly damped ion waves, and nonlinear ion heating can decrease the total SBS growth. In the opposite limit $\ell/L > 1$ the total number of exponentiations will be independent of ν_i . Ion heating is then not predicted to cause significantly decreased SBS growth. Under circumstances where inhomogeneity dominates, the main effect of ion heating could in some cases be to bring SBS down to its collisional threshold.

3.1.3 Effect of Light Reflected from the Critical Surface

The relation between reflectivity R and quality factor Q_H given in Eq. (43) can be changed dramatically. Light which has been classically reflected from the critical-density surface has the correct frequency to play the role of the back-reflected wave in SBS. This can be the case if the plasma outflow speed relative to the critical density contour is not supersonic locally. Under these conditions the light reflected from critical acts as a very large noise source for the backscattered wave [51], and thus for the same Q_H the scattering can be more intense than it would be in the absence of the classically reflected light.

To see this effect quantitatively [51], consider a plasma with outflow velocity $v(x)$ and density $n(x)$ whose critical-density contour is moving with speed v_c . [Thus the speed of the plasma relative to the n_c contour at $x = x_c = v(x_c) - v_c$.] Consider the implications of the matching conditions. In the rest-frame of the plasma, the frequency of the incident laser light as seen by the plasma is $\omega_i = \omega_1 + k_1 v(x)$, where ω_1 and k_1 are the laser light wave's frequency and vacuum wavenumber. The frequency of light classically reflected from the n_c surface is $\omega_R = \omega_1 - k_1 [v(x) - 2v_c]$, because of the Doppler-shift at the moving critical surface. From wavenumber matching, the wavenumber of the SBS ion-acoustic wave at a density $n_c(x)$ must be $k_s = 2k_1 \epsilon_c [n(x)]^{1/2}$, where $\epsilon(x) = 1 - n_c(x)/n_0$ is the local dielectric constant for the incident and reflected light waves. Thus the frequency-matching requirement $\omega_{it} = \omega_R + \omega_{sa}$ is satisfied exactly at the point x_* where $v(x_*) = v_c + v_c [\epsilon(x_*)]^{1/2}$. At this position in the plasma, the classically reflected light can provide an enhanced noise source to drive strong SBS. The resonance condition

$$\frac{v(x_*) - v_c}{v_c} = [\epsilon(x_*)]^{1/2} \approx 1 \quad (48)$$

thus requires that the outflow speed *relative* to the critical-density contour be subsonic. (Note that if the critical density surface is itself moving rapidly outward, as it tends to do during the rising part of the laser pulse, this does not necessarily mean the plasma is subsonic in the *lab* frame.)

To find the reflectivity when this resonance condition is satisfied, we solve Eqs. (39) and (42) for the simple example of a homogeneous plasma slab moving uniformly with the speed given by Eq. (48), and bounded on one end by a totally reflecting critical surface [51]. The geometry is shown in Fig. 15. At $x = 1$ the laser is incident with amplitude $A_i(x = 1) = A_1$. When the left-going (or "incident") wave reaches $x = 0$, its amplitude has fallen to $A_i(x = 0) = T^{-2} A_1$, where T is the transmission and will be found shortly. Because of the perfectly reflecting critical surface just to the left of $x = 0$, the back-reflected amplitude A_R is also equal to $T^{-2} A_1$ at $x = 0$.

FIG. 15. Boundary conditions for calculation of SRS including light reflected from critical surface [51]. The transmission coefficient is T . The outward speed of the uniform plasma relative to the critical density contour is $v - v_c$.

Equations (39) and (42) may be combined to yield

$$\frac{d}{dx} A_i^2 = \frac{d}{dx} A_b^2 = - \left(\frac{Q_H L}{A_i^2} \right) A_i^2 A_b^2, \quad (49)$$

where we have used the definition of Q_H in Eq. (44) for convenience of notation. The solution of the first equation in (49) is $A_i^2 = A_b^2 + \text{const}$. This can be substituted into the second equation in (49) and integrated to obtain

$$T = \frac{1}{1 + Q_H} \quad (50)$$

Figure 16 shows this relation between the transmission and Q_H , and compares it with the prediction for partial reflection from n_c and for no reflection from n_c [Eq. (43)]. Even with only 25% reflection from the critical surface, the transmission falls much more quickly at small Q_H than it does if the critical-surface effect is neglected. This is because the back-reflected noise signal no longer has to exponentiate 10 or 20 times to become comparable with the incident intensity. In a more realistic situation where the resonance condition (48) is not exactly satisfied, the effect seen in Fig. 16 persists when the ion damping rate $\nu_i \propto \omega_{pi}$ is not too small [51].

There are several questions about SRS with light reflected from n_c which still need to be addressed. First, the above analysis was for a one-dimensional situation. For oblique incidence on a density gradient, light is specularly reflected from n_c and is subject to refraction along its ray path. The physics of the resonance condition will thus change. Second, if the critical-density surface is rippled, light from n_c will be diffusely reflected into a wide range of angles. Thus even for normal incidence, the fraction back-reflected from n_c which can feed the above resonance mechanism may be as low as 5×10^{-7} . Third, the possibility of absolute instability arises in this situation, both because of the required subsonic flows and because of feedback via

FIG. 16. SBS transmission coefficient [51] T , versus quality factor Q_H [Eq. (44)]. The curve labeled "0% reflection" represents SBS which grows only from thermal noise, with noise level $I_{\omega} I_{\omega} = 10^{-4}$. The other two curves show SBS transmission coefficients including the effect of light reflected from the critical surface, with reflection fractions from the critical surface of 27% and 100%, respectively. The curve labeled "100%" is given by Eq. (50). All three of these curves are for exact resonance, Eq. (48).

reflection from the critical density surface. Finally, Randall [52] has seen in the fluid computer code AURUS that profiles of plasma density and velocity can self-adjust in the presence of SBS fed by reflection from the critical surface, so that the strong reflectivity occurs in a relatively localized part of the profile which then moves in the course of time. Such phenomena might lead to a more dynamical picture, in which reflection-driven SBS occurs in different parts of the underdense plasma at different times in the laser pulse, as the profiles self-adjust to the scattering.

3.1.4 Competition Between SBS and Inverse Bremsstrahlung Absorption

One of the most straightforward ways to limit SBS is to use laser and target parameters where inverse Bremsstrahlung is strong. When absorption is rapid enough, the distance over which SBS can grow is limited to a few times the absorption length for inverse Bremsstrahlung, κ_{ib}^{-1} . Thus if the growth length for SBS is longer than a few times κ_{ib}^{-1} , the light will be absorbed before it can be scattered by SBS. This limitation mechanism is particularly effective for short-wavelength lasers, where classical absorption tends to be large.

In a uniform plasma, the outcome of the competition between SBS and inverse Bremsstrahlung is sensitive to the plasma's electron density [53]. This is because the growth rate of SBS is proportional to n_e [Eq. (44)], whereas the classical absorption length κ_{ib}^{-1} is proportional to n_e^2 [Eq. (16b)]. Thus SBS always dominates if the density is low enough, and absorption wins if the density is high enough.

However in a laser-produced plasma with a gradient in density from $n_c = 0$ to $n_c = n_e$, the mix between low-density and high-density regions (and hence between SBS and classical absorption) is determined by the characteristics of the profile. We can illustrate this effect by the following simple example [54]. Consider strong inverse Bremsstrahlung in a spherical expansion with density profile

$$n_e(r) = n_e \left(\frac{r_0}{r} \right)^\beta, \quad \beta > 1. \quad (51)$$

and with constant coronal temperature. In the absence of Brillouin scattering, the radius r_{ib} at which inverse Bremsstrahlung becomes important can be calculated from Eq. (17) by setting $i\lambda_{ib}(r)dr = 1$ and using the density profile (51). Then in the presence of SBS, we can roughly parametrize the Brillouin reflectivity by an analogue of the quality factor Q_{ib} discussed in Section 3.1.2 above:

$$Q_{ib} = \frac{\pi}{2} \left(\frac{E_L^2}{8\pi n_e T_e} \right) \int_{r=z}^{r_{ib}} \frac{n(r)}{n_e} \frac{dr}{\lambda_L} \quad (52)$$

which is the number of SBS gain lengths in one absorption length.

With the density profile given in (51), this quality factor for SBS in the presence of inverse Bremsstrahlung scales as

$$Q_{ib} \propto \left(\frac{E_L^2}{8\pi n_e T_e} \right) \left(\frac{r_0}{\lambda_L} \right) \left[\frac{r_0(n_e)}{c} \right]^{\frac{1-\beta}{2\beta-1}} \quad (53)$$

SBS can be strongly limited by inverse Bremsstrahlung provided that $Q_{ib} \leq 5$ or so. Thus the result (53) shows clearly the advantages of short laser wavelengths λ_L in limiting SBS by this mechanism, since

$$Q_{ib} \propto \lambda_L^{\frac{4\beta-1}{2\beta-1}} \quad (54)$$

where the exponent of λ_L lies between 1 and 2 for $\beta > 1$. For a simple linear profile $n_e(r) = n_e \times (r/r_0)$, the quality factor Q_{ib} is reduced even more strongly for short laser wavelength [55]: $Q_{ib} \propto \lambda_L^{2/3}$.

To study the competition between SBS and inverse Bremsstrahlung more carefully, the WKB propagation of the incident and Brillouin-scattered electromagnetic waves was integrated numerically in Ref. [54]. A steady-state flow in spherical geometry was taken for the background plasma conditions. Both incident and scattered light underwent inverse Bremsstrahlung absorption, and 30% of the light reaching the critical surface was absorbed there to model resonance absorption. The (isothermal) electron temperature was taken as given, but the ion temperature profile was obtained from an ion energy equation which included collisional coupling to electrons, expansion cooling, and a simple main-body model for nonlinear ion heating due to SBS [46]. A representative result can be illustrated by using the conditions described in Fig. 17, which might be characteristic of large, reactor-sized targets (critical radius = 0.24 cm, laser power $P_L = 145 \times 10^{12}$ W, $T_e = 1.5$ keV in the corona).

The resulting steady-state absorption and Brillouin backscattering are shown in Fig. 18 for three different laser wavelengths: $\lambda_L = 0.3 \mu\text{m}$, $0.5 \mu\text{m}$, and $1.0 \mu\text{m}$. If there were no SBS or inverse Bremsstrahlung, the incident wave amplitude would be expected to swell as it propagates from large radii toward the critical radius $r_c = 0.24$ cm, both due to spherical convergence ($r^2 E^2 = \text{const.}$) and due to the decrease in the group velocity near n_c . The parameters for Fig. 18a, $\lambda_L = 1 \mu\text{m}$, were chosen to illustrate the behavior of the incident and Brillouin-scattered waves when there is no SBS stabilization by inverse Bremsstrahlung. The natural swelling described above is completely overcome by strong SBS at low densities. As the incident wave approaches r_c its amplitude does increase somewhat, but because the wave has already been depleted by SBS only 4% of the laser light energy is absorbed.

Figures 18b and 18c show the stabilization of SBS that occurs at shorter laser wavelengths. For $\lambda_L = 0.5 \mu\text{m}$ (Fig. 18b), one is just entering the regime where inverse Bremsstrahlung becomes significant. There is still some SBS which depletes the incident beam at low density, but inverse Bremsstrahlung is sufficiently strong as the wave approaches n_c that the total absorption has risen to 25%.

At the shortest laser wavelength $\lambda_L = 0.3 \mu\text{m}$ (Fig. 18c), inverse Bremsstrahlung stabilization is very strong. The low level of SBS allows the incident wave to swell freely as it propagates toward smaller radii. At

FIG. 17. Density and velocity profiles used for numerical study [54] of SBS, including effects of inverse Bremsstrahlung absorption. Spherical target of critical radius $r_c = 0.24$ cm was irradiated with laser power $P_l = 1.45 \times 10^{14}$ W. Coronal electron temperature was 1.5 keV

high densities inverse Bremsstrahlung is sufficiently potent to overcome the swelling and actually deplete the incident wave energy before it reaches critical. As a consequence 93% of the laser energy is absorbed.

This type of numerical work confirms the potential importance of inverse Bremsstrahlung for the stabilization of Brillouin at laser wavelengths less than one micron.

The results of the above discussion of SBS physics lead to the following overall conclusions. As of the present time, there seem to be several ways to reduce Brillouin scattering:

- *Ion Wavebreaking.* At the highest laser intensities and longest laser wavelengths, ion wavebreaking or trapping can limit the ion-wave amplitude to 10–20%. However if the length of the underdense plasma is several hundred vacuum wavelengths or more, SBS can be large nevertheless.

FIG. 18. Electric field profiles for incident (solid line) and backscattered (dashed line) light waves, including effects of SBS and inverse Bremsstrahlung absorption [54]. Plasma conditions are described in Fig. 17. Profiles are shown for laser wavelengths a) $\lambda_L = 1.0 \mu\text{m}$ (4% absorption), b) $\lambda_L = 0.5 \mu\text{m}$ (25% absorption), and c) $\lambda_L = 0.3 \mu\text{m}$ (93% absorption). SBS decreases as the laser wavelength becomes shorter.

• **Nonlinear Ion Heating.** This is most important for laser intensities and wavelengths just below the wavebreaking limit. By raising the damping rate of the ion wave, ion tail heating can decrease the total SBS gain in a nearly uniform plasma. The effect is weaker in a plasma with strong gradients.

• **Inverse Bremsstrahlung.** In the complementary regime of shorter laser wavelengths and lower laser intensities, strong inverse Bremsstrahlung absorption can limit the growth distance for SBS. In this case the geometrical length L used to find the SBS exponentiation is replaced by the (shorter) absorption length L_{abs} .

The following two additional stabilization mechanisms have not been discussed in the present report, but they are thoroughly described in Refs. [56] and [57]:

• **Broad Laser Bandwidth.** If the spectral width of the laser light is larger than the effective SBS gain width, the SBS growth can be significantly decreased.

• **Irregularities in the Background Plasma.** For reactor-sized targets SBS is predicted to be large not because the Brillouin spatial growth rate is huge, but rather because there are 10^2 – 10^3 vacuum wavelengths available to grow in. Under these conditions a process which destroys the coherence of the incident, backscattered, and ion-acoustic waves can thus diminish the convective amplification considerably, replacing the geometrical length available for growth by the (shorter) coherence length. It can also raise the instability thresholds accordingly.

3.2 STIMULATED RAMAN SCATTERING AND TWO-PLASMON DECAY INSTABILITY

In stimulated Raman scattering [38,40,48,53,58,59], the incident laser light wave decays into an electron plasma oscillation and a scattered electromagnetic wave. Since the frequency of the plasma wave (ω_{pw}) is $\omega_{\text{pw}} = \omega_{\text{pe}}(1 + 3k_{\text{pw}}^2 \lambda_{\text{De}}^2)^{1/2}$, the wavenumber and frequency-matching conditions can only be satisfied for densities $n_e < n_c/4$. The linear growth rate maximizes for backscatter, and is equal to

$$\gamma_{\text{Raman}} = (k_{\text{pw}} v_{\text{th}} / 2\sqrt{2}) [\omega_{\text{L}} - \omega_{\text{pe}}]^{1/2},$$

where $v_{\text{th}} \equiv eE_{\text{L}} / m_e \omega_{\text{L}}$.

At the quarter-critical surface $n_e = n_c/4$, a second process called the two-plasmon decay instability [60], or “ $2\omega_{\text{pe}}$ instability,” can compete with Raman. In this case the incident light wave decays into two Langmuir waves or plasma oscillations, and for $k \gg k_L$ the wavenumber matching conditions then require the Langmuir waves to propagate almost opposite each other at angles near 45° to the incident $\mathbf{k}_L$ of the laser. When $k \approx k_L$, Landau damping causes the k s of the two plasma waves to become unequal and shifts their relative angles away from 45° . This is the case for conditions appropriate to many laser fusion targets. The maximum growth rate is $\gamma_{2\omega_{\text{pe}}} \approx \omega_L v_{\text{th}}/4c$. The frequency and wavenumber matching conditions for Raman and $2\omega_{\text{pe}}$ decay are illustrated in Fig. 19.

The behavior of the Raman and $2\omega_{\text{pe}}$ processes can be divided into two different regimes, depending on the value of the local density n_e/n_c . In the region near one-quarter of the critical density, $0.2 < n_e/n_c \leq 0.25$, Raman and $2\omega_{\text{pe}}$ occur together. Raman backscatter in particular is an absolute instability at $n_e = n_c/4$, because the backscattered light wave has zero group velocity there [61]. The $2\omega_{\text{pe}}$ instability can also be absolute. In addition, the $2\omega_{\text{pe}}$ and Raman thresholds due to inhomogeneity are relatively low [50] near $n_c/4$.

As a consequence of the low threshold and absolute nature of the instabilities, the nonlinear behavior of Raman and $2\omega_{\text{pe}}$ near $n_e = n_c/4$ is very strong and contains a wide variety of different physical processes. Particle-in-cell computer simulations at high laser intensity [62] show that in this regime the instability seems to operate in bursts. The rough chronology is as follows [62]: If one begins with a relatively gentle density gradient near $n_c/4$, the Raman and $2\omega_{\text{pe}}$ instabilities are mixed together and grow rapidly from noise. When they reach large amplitude, the ponderomotive force ($E^2/8\pi$) due to all the different plasma waves is strong enough to drive large ion density fluctuations. At the same time a heated electron tail is generated [48]. In Ref. [62] the instabilities were damped primarily by the strong ion density fluctuations,

Raman

$$\left. \begin{aligned} \vec{k}_L &= \vec{k}_R + \vec{k}_{pw} \\ \omega_L &= \omega_R + \omega_{pw} \end{aligned} \right\} \begin{aligned} \omega_{pw} &= \omega_p \left(1 + 3 k_{pw}^2 \lambda_{De}^2 \right)^{1/2} \\ |k_{R,L}| &= \frac{\omega_{R,L}}{c} \left(1 - \frac{n_e}{n_c} \right)^{1/2} \end{aligned}$$

Two-Plasmon Decay

$$\left. \begin{aligned} \vec{k}_L &= \vec{k}_{p1} + \vec{k}_{p2} \\ \omega_L &= \omega_{p1} + \omega_{p2} \end{aligned} \right\} \begin{aligned} k_L &= \frac{\omega_L}{c} \left(1 - \frac{n_e}{n_c} \right)^{1/2} \approx \frac{\sqrt{3}}{2} \frac{\omega_L}{c} \\ \omega_{p1,p2} &= \omega_p \left(1 + 3 k_{p1,p2}^2 \lambda_{De}^2 \right)^{1/2} \end{aligned}$$

FIG. 19. Wavenumber and frequency matching conditions for the Raman and two-plasmon decay instabilities.

which locally steepened the profile sufficiently to push the instabilities below their inhomogeneity thresholds. Figure 20 illustrates the steepened profile.

Once the profile is steepened and the hot electrons are generated, the instabilities can no longer continue to grow. The ion waves and hot electrons gradually dissipate, the profile relaxes, and conditions are appropriate for a second burst of instability growth.

Reference [63] discusses another nonlinear mechanism for the $2\omega_{pe}$ instability at $n_e/4$, which has not been observed in computer simulations at high or moderate laser intensities.

Heated electrons from the instabilities at $n_e/4$ are a major preheat concern. For a high-intensity, steep-profile irradiation such as that of Ref. [62], the time-averaged fraction of the laser light energy in the heated electron tail is $<10\%$; the temperature of these tail electrons can be quite high: $T_h \approx 75$ keV for a plasma with background temperature of 10 keV and an incident laser intensity of 10^{16} W/cm², $\lambda_l = 1.06$ μ m [62]. Future targets, such as those intended for eventual reactor lasers, are expected to have underdense plasmas that have longer spatial extents in the neighborhood of $n_e/4$. Under these circumstances, the instabilities at $n_e/4$ are capable of putting as much as 30–40% of the laser energy into hot-electron tails. These tails can be extremely hot; for example, Ref. [64] found $T_h = 110$ keV for $I_l = 3 \times 10^{15}$ W/cm², $\lambda_l = 1.06$ μ m, $T_e = 3$ keV $= 3 T_i$, for a flat plasma slab 21 vacuum wavelengths long at $n_e = 0.23 n_c$.

The second interesting regime for Raman scattering [64] occurs at intermediate densities where $n_e < 0.2 n_c$, but n_e is not so small that the plasma wave is heavily damped. (Strong damping takes place when $k_{pw} \lambda_{De} > 1/3$.) In this intermediate-density regime the $2\omega_{pe}$ processes cannot occur, and Raman is convective except for exact sidescatter [65]. Thus Raman is less violent than it is near $n_e = n_c/4$. In addition the intensity threshold due to inhomogeneity is higher than it is at the quarter-critical surface.

In this regime the nonlinear behavior is less complex, and computer simulations at high laser intensity indicate [64] that self-consistent electron tail heating can limit Raman scatter. The argument is quite analogous to that described earlier for ion tail heating in Brillouin scattering. Figure 21 shows a typical heated electron distribution, which has a nearly Maxwellian tail of hot electrons with characteristic temperature

FIG. 20. Steepened electron density profile near the quarter-critical density surface, due to the two-plasmon decay and Raman instabilities [62]. Computer simulation parameters are $v_a/c = 0.1$, initial $T_e = 1$ keV, $\lambda_{I_i}/m_e = 100$, $ZT_e/T_i = 1$.

$$T_h \approx 64 \text{ keV} \left(\frac{(n_e/n_c) (1 + 3k_{pw}^2 \lambda_{De}^2)}{(1 - n_e/n_c) (1 - 1/2 \sqrt{n_e/n_c})^2} \right) \quad (55)$$

$$n_e \ll n_c, \quad k_{pw} \lambda_{De} < \frac{1}{3}.$$

The heated electron temperature is appropriate to provide damping for the electron plasma wave.

The consequences of electron tail heating in Raman [65] are similar to those of ion tail heating in Brillouin. For a homogeneous medium, the Raman quality factor (analogous to Eq. (44) for SBS) is

$$Q_H \approx \left(\frac{k_{pw}^2 L}{8k_R} \right) \left(\frac{eE_L}{m_e \omega_L c} \right)^2 \left(\frac{\nu_{pw}}{\omega_{pe}} \right)^{-1}, \quad (56)$$

where k_{pw} and ν_{pw} are the wavenumber and damping rate of the plasma wave. In a homogeneous medium the heated electrons decrease Q_H by increasing the plasma wave damping ν_{pw} . Then by the Raman analogues of Eq. (43) and Fig. 13, the reflectivity R is decreased as well. A consequence of the lower reflectivity is a population of hot electrons with $50 \text{ keV} < T_h \leq 100 \text{ keV}$. For a strongly inhomogeneous medium well above the

FIG. 21. Electron distribution function produced by Raman scattering, from computer simulation by K. Estabrook. The heated electron tail is evident. This 1-D simulation used an incident laser intensity of 10^{16} W/cm 2 , $\lambda_L = 1.06$ μ m, initial $T_e = 10$ keV, and density profile rising from $n_e = 0$ to $n_e = 0.41 n_c$ in a distance of $70 \lambda_L$.

damping threshold, increased plasma-wave damping due to electron tail heating should not affect Raman as strongly [49]. This is analogous to the result for SBS.

As was the case with SBS, Raman scatter may additionally be limited by inverse Bremsstrahlung absorption or broad laser bandwidth. We discuss these in turn.

The absolute instabilities at $n_e = n_c/4$ occur over a very narrow range of density. Thus the main way by which inverse Bremsstrahlung could affect them is by diminishing the laser intensity that actually reaches the $n_c/4$ surface. If the local intensity is diminished enough, Raman and $2\omega_{pe}$ could be pushed below threshold. Numerically this process is not important until [64,65] the laser wavelength is a fraction of a micron, or else until the laser intensity is $\leq 10^{13}$ W/cm 2 .

The physics of inverse Bremsstrahlung stabilization of Raman is somewhat different in the convective regime $n_e < 0.2 n_c$. The process has not yet been examined in detail. But because inverse Bremsstrahlung falls off at low densities as $(n_e/n_c)^2$, it seems numerically more difficult to find a situation where enough light could be classically absorbed in the very underdense plasma to seriously affect Raman scattering there. (In the comparable situation for Brillouin the convective growth of SBS continues all the way to $n_e \approx n_c$, so that strong inverse Bremsstrahlung absorption between $n_e \approx n_c/4$ and $n_e \approx n_c$ can diminish SBS substantially.)

As was the case with SBS, broad laser bandwidth may also reduce the growth of Raman scattering [56,57,64]. For a finite plasma of length l , the laser bandwidth $\Delta\omega_L$ must satisfy two inequalities [56] in order to decrease Raman growth:

$$\frac{\Delta\omega_L}{\omega_L} > \frac{\gamma_0^2}{\omega_L^2} \left\{ \omega_L l/c \right. \quad (57)$$

$$\left. \omega_L/\nu_{pw} \right.$$

Here ν_{pa} is the plasma wave damping rate and γ_0 is the Raman temporal growth rate. Typically these inequalities mean that the minimum required bandwidth to stabilize Raman is larger than that for SBS. Computer simulations [65] indicate that such wide-bandwidth or multiline lasers can effectively reduce Raman scatter in some parameter regimes, particularly for laser intensities less than a factor of 10 above threshold.

Competition between Raman sidescatter and backscatter was assessed in Ref. [66], for $n_e < n_c/4$. Recent experiments at the Lawrence Livermore National Laboratory indicate that Raman sidescatter can be significant.

3.3 FILAMENTATION AND SELF-FOCUSING

Filamentation and self-focusing refer to the process by which an intensity modulation initially present on the laser beam can become more pronounced as the light propagates through the plasma toward the critical density surface [38,67]. If the initial intensity modulation involves only a small part of the laser beam, the process is generally called filamentation. Self-focusing describes the mechanism when it involves the large-scale intensity profile of the laser beam as a whole.

The physics of filamentation or self-focusing at high laser intensity is illustrated schematically in Fig. 22. A region of initially higher $E^2/8\pi$ pushes plasma aside because of its additional light pressure. The consequent lower plasma density n_e increases the index of refraction of the plasma

$$n = \frac{k_1 c}{\omega_1} = \left(1 - \frac{n_e}{n_c}\right)^{1/2} \quad (58)$$

in the high-intensity region, thus bending the wave front as shown in Fig. 22. Since wave energy flux propagates perpendicular to the wave front, the curvature of the wave front produces a focusing effect which increases the light intensity still more.

This is an unstable process which convectively amplifies any intensity inhomogeneity initially present on the laser beam. Since high-power laser beams are generally far from uniform (see Fig. 23), they should in principle provide ample initial inhomogeneities for the plasma to amplify.

Theories for the nonlinear consequences of filamentation fall into two contrasting categories: 1) The filament forms and is able to remain more or less intact while it propagates to the critical density surface. 2) The filament is not able to retain its shape over substantial distances after it forms, due to the disruptive effects of stimulated scattering, diffraction, or refraction in the plasma. We will discuss these alternatives in turn, and try to indicate their consequences. It is not yet clear which of the two scenarios is more likely to occur.

3.3.1 Filaments Stay Intact and Propagate to the Critical Surface

If this is the case, filamentation could have major implications for processes such as electron heating and absorption at the critical surface, and for the structure of laser-produced dc magnetic fields.

FIG. 22. Physical mechanism for filamentation driven by the ponderomotive force.

FIG. 23. Spatial profiles of the Argus laser focal spot at three different incident laser powers ($1 \text{ TW} = 10^{12} \text{ W}$). These scans, taken at the Lawrence Livermore National Laboratory in 1977, show substantial spatial modulation at terawatt laser powers. In the years since 1977, beam relaying techniques developed by J. Hunt and colleagues have resulted in considerably better beam homogeneity. (Data shown here are from W. Simmons, D. R. Speck, and J. Hunt.)

Kruer has recently proposed a simple model for hot-electron generation in a filament that reaches n_c . Or it way to the critical surface, self-focusing can cause the intensity inside the filament to increase until roughly $\approx E_L / m_e \omega_L$ inside the filament becomes comparable to v_{te} in the surrounding plasma. If we suppose we are in the high-intensity regime where hot electrons dominate the underdense regions, Eq. (35) implies that

$$\left(\frac{I_L}{10^{15} \text{ W/cm}^2} \right) \left(\frac{\lambda_L}{1.06 \mu\text{m}} \right)^2 \approx 2.2 \left(\frac{T_h}{1 \text{ keV}} \right) \quad (59)$$

On the other hand once the filament finally reaches the n_c surface, it will make hot electrons there at a temperature T_h given by Eq. (30) or (31). We can combine Eq. (59) with one of these two equations and solve uniquely for the intensity inside the filament and for T_h , assuming $T_e \approx 1 \text{ keV}$. If we choose Eq. (30), we obtain

$$\left(\frac{I_L}{10^{17} \text{ W/cm}^2} \right) \left(\frac{\lambda_L}{1.06 \mu\text{m}} \right)^2 \approx 2.5 \cdot \left(\frac{T_h}{1 \text{ keV}} \right) \approx 110, \quad (60a)$$

whereas Eq. (31) yields

$$\left(\frac{I_L}{10^{17} \text{ W/cm}^2} \right) \left(\frac{\lambda_L}{1.06 \mu\text{m}} \right)^2 \approx 0.6 \cdot \left(\frac{T_h}{1 \text{ keV}} \right) \approx 30 \quad (60b)$$

In both cases the final intensity and hot electron temperature inside the filament are quite large.

A second consequence of filaments which stay intact and propagate up to n_c is increased density-profile deformation and resonance absorption at the critical surface. This is seen clearly in a particle-in-cell computer simulation by Langdon [68], in which the laser beam was normally incident and was initially spatially modulated by a factor of 4 across the incident beam profile. The filament which resulted from the initial modulation propagated toward the critical surface. When it reached n_c , the increased light pressure inside the filament made a local indentation in the surface, as shown in Fig. 24. On the sides of the indentation the local angle between ∇n and the laser's E field became appropriate for strong resonance absorption, despite the fact that the laser beam had originally been at normal incidence. As a result of the deformation of the n_c surface produced by the filament, resonance absorption was locally increased.

A third possible consequence of filaments which remain intact is the production of dc magnetic fields around the outside of the filament [69]. We shall show in Section 5.2 that there is a source for dc magnetic fields in a plasma which is proportional to $(e/c) \nabla n \cdot n_c \times \nabla T_e$. In an underdense plasma with the geometry shown in Fig. 25, this source would produce B fields which encircle the filament. Such channel-like structures could explain the localized hot-electron emission and the "filamentary" x-ray patterns seen by Donaldson and Spalding [70], since if the B fields were strong enough they could confine hot electrons to the interior of the filament. These B fields surrounding light filaments might also explain the filament-like magnetic structures seen by Willi and Rumshy [70] in long-pulse experiments.

To date, most formal analytic theory for the nonlinear evolution of filaments [71] has concentrated on the case where they remain intact. Ion inertia is generally neglected, so that the local density is given by the quasistatic relation

$$n_e \approx n_0 \exp \left(\frac{-e^2 \langle E^2 \rangle}{2m_e \omega_L^2 T_e} \right) \quad (61)$$

The background plasma density n_0 is taken to be uniform. It is also frequently assumed that the transverse intensity distribution follows a Gaussian shape, whose width and height can vary with distance along the filament.

The main results of Ref. [71] are in qualitative agreement. (A) For any incident laser power P_L , there exists a particular filament radius, called the "self-trapping" radius. A filament with this initial radius can propagate in a steady state, without additional spreading or focusing. Under these conditions the nonlinear index of refraction [which results when the density distribution (61) is substituted into (58)] just balances diffraction. (B) If the initial filament radius is not equal to this "self-trapped" value, propagation of the filament is

FIG. 24. Indentation of the critical-density surface, caused by spatial modulation of a factor of 4 across the incident laser beam profile. This figure shows isodensity contours from a particle-in-cell computer simulation by A. B. Langdon [68]. Simulation parameters were $I_1 = 10^{16} \text{ W cm}^{-2}$, $\lambda_1 = 1.06 \mu\text{m}$, initial $T_e = 5 \text{ keV}$, $M_i/m_e = 100$, $T_e/T_i = 3$, $\omega_L t = 700$.

oscillatory, as shown in Fig. 26. The filament begins to come to a strong focus due to the ponderomotive nonlinearity. But the density deformation given in (61) saturates when $e^2 E^2 / 2m_e \omega_L^2 T_e$ is large, since the density cannot decrease below zero. Thus near a strong focus, diffraction can become larger than nonlinear self-focusing, and the filament spreads out in radius again. If the filament is assumed to always maintain a nearly Gaussian intensity profile, this process of inward focusing and outward diffraction is predicted to repeat indefinitely.

3.3.2 Filaments Fall Apart Before They Propagate Very Far

The second scenario for the nonlinear evolution of filaments implies that they cannot stably propagate for long distances in the plasma.

There are several possible reasons for lack of filament stability. Several years ago, Valen [72] showed that filaments in two dimensions can be unstable to bending perturbations, where an initially straight filament becomes kinked and distorted. The geometry is shown in Fig. 27a. If this kink or bending instability is found to persist in three dimensions with the full exponential density nonlinearity (61), it would be difficult for a filament to propagate for long distances.

When the background plasma density is above $n_c = n_c/4$, a quarter-critical density surface can be formed on the walls of a filament, because of the ponderomotive density depression at the filament's center. When this occurs, Langdon and Lasinski [73] showed in computer simulations that a vigorous form of the two-plasmon decay instability can produce electron acceleration inside the filament. Figure 27b shows the strongly modulated electrostatic potential which results from the $2\omega_{pe}$ decay process inside the filament. The

FIG. 25. Mechanism for production of $\nabla n_e \times \nabla T_e$ magnetic field surrounding a light filament [69].

two-plasmon decay and electron acceleration act as an energy drain on the light inside the filament, thus depleting the internal high-intensity region.

More recently, nonlinear phenomena which accompany the "repeated foci" shown in Fig. 26 have begun to be investigated. Lasinski and Langdon [74] have followed numerically the evolution of a focusing filament without making a Gaussian-shape hypothesis for the intensity distribution. They found that while the central part of the beam underwent the predicted repeated foci [71], the outer part of the beam did not participate. Figure 27c shows that only the inner part of the beam remains to propagate intact in the manner of Fig. 26.

In related work, Randall [74] has considered the enhancement in Brillouin backscatter which occurs when a filament begins to focus to high intensities. He finds for some experimental conditions that during the focusing process, much of the light in the filament can be Brillouin backscattered before the focus itself is reached. This is illustrated in Fig. 27d. Thus the work of Lasinski and Randall seems to imply that a "real" filament would not be likely to survive more than one or two repeated foci, both due to SBS and due to non-homologous behavior of the exterior parts of the beam.

As the reader will have gathered from the above discussion, theories for the eventual fate of filamentation are not yet definitive. In particular, nonlinear phenomena occurring near a strong focus have only recently begun to be explored. It seems to me that this issue of the physics near strong foci is crucial to understand, because only with very special initial conditions can a filament propagate in a "self-trapped" manner [71] without beginning to focus.

For very low laser intensities and collisional plasmas, filamentation driven by the Joule heating non-linearity can sometimes become competitive with the ponderomotive force-driven mechanism. This process is reviewed by Sodha et al. [71], and was discussed in Ref. [75] for an application to the earth's ionosphere.

FIG. 26. Schematic picture of "repeated foci" predicted by nonlinear theories of self-focusing [71].

FIG. 27. Four ways in which the integrity of a light filament might be destroyed before it reaches the critical surface. In all four cases the laser beam is incident from the left. a) Bending instability [72]. b) Two-plasmon decay instability inside a filament [73]. The figure shows contours of electrostatic potential ϕ . The modulation is due to the two daughter-wave plasmons produced by the decay instability. The units of distance are $2c/\omega_L$. c) Lack of participation of the outer parts of the beam in the self-focusing process, from a calculation by Lasinski and Langdon [74]. These flux contours would have remained equally spaced if all of the initially Gaussian beam had self-focused together. d) Effect of Brillouin scatter near a nonlinear focus, from calculation by Randall [74]. The mean radius "r" of the filament diverges indefinitely after the first focus, because strong SBS has depleted the beam intensity and hence weakened the self-focusing nonlinearity.

4. INHIBITION OF ELECTRON THERMAL CONDUCTION

4.1 INTRODUCTION AND BASIC CONCEPTS

An essential feature of laser fusion targets is that laser energy absorbed in the corona must be effectively conducted toward colder, high-density regions of the target. Figure 28 illustrates a typical temperature profile for a thick initially-solid target. A large inward heat flux from the critical surface towards high-density regions means that newly heated material at the surface of the cold dense target will expand outward rapidly. Then by conservation of momentum the cold remainder of the target gets a bigger inward push due to the rocket effect. Figure 29 illustrates this process. Conversely, inhibited heat flux leads to slower outward expansion of the blow-off. The inner parts of the target then implode less vigorously.

When the mean free path λ_{mfp} is small compared to the temperature gradient scale length L_T , the classical heat flux in a plasma can be written as

$$\mathbf{q} = -\kappa \nabla T \text{ ergs/cm}^2 \text{ s} \quad (62)$$

where to order of magnitude

$$\kappa \nabla T \approx \left(\frac{\lambda_{mfp}}{L_T} \right) (nT) v_t, \quad v_t \approx \sqrt{T/M} \quad (63)$$

If electrons and ions have equal temperatures and densities, the first two factors $(\lambda_{mfp}/L_T)(nT)$ are the same for both species. However because the electron thermal speed v_{te} is larger than the ion thermal speed by the factor $(M_i/m_e)^{1/2}$, it is the *electron* thermal conductivity which dominates in Eq. (63). Thus in a plasma, it is the random walk of the electrons which carries heat from hot regions to cold regions.

Classical transport coefficients are calculated assuming the plasma is sufficiently collisional that $\lambda_{mfp} \ll L_n, L_T$, where $L_n \equiv n^{-1} |dn/dx|$ and $L_T \equiv T/|dT/dx|$, and assuming that $\nu_{ei} \gg \omega_{ci}$. Under these conditions the electron distribution function is locally Maxwellian, and the electron heat flux is [10,19]

$$\mathbf{q}_e = -\kappa \nabla T_e \text{ ergs/cm}^2 \text{ s} \quad (64)$$

FIG. 28. Typical form of electron temperature profile in a laser-heated initially solid target. The temperature is high in the corona where laser absorption occurs, and cold in the high-density regions. Thus an inward heat flux flows from $n_e = n_c$ toward higher densities.

FIG. 29. Inward electron heat flow produces an outward flow of newly heated plasma, in a spherically illuminated target. In the presence of this hot plasma outflow, the dense part of the target recoils inward to conserve momentum. The amount of inward heat flux plays a crucial role in determining the magnitude of the rocket-like recoil and implosion.

$$\kappa_k = 20 \left(\frac{2}{\pi} \right)^{3/2} (\epsilon \delta_T) \left[\frac{(kT_e)^{5/2}}{(m_e)^{1/2} c^4 Z \ln \Lambda} \right] \quad (65)$$

$$\epsilon \delta_T \approx \frac{0.095 (Z + 0.24)}{1 + 0.24Z} \quad (\text{to } \pm 10\%) \quad (66)$$

The factor $\ln \Lambda$ is the usual Coulomb logarithm. Note that κ depends on density only weakly, through $\ln \Lambda$.

In an underdense plasma, these classical transport coefficients can be modified by the presence of intense laser light. Since the Coulomb collision rate is proportional to v^{-3} , the oscillating velocity $v_0 = eE_1 / m_e \omega_L$ of an electron in the laser electric field E_1 can modify collisional transport when v_0 becomes comparable with the electron thermal speed v_{te} . This subject has been discussed in detail in Ref. [76].

In general the population of "hot" or suprathermal electrons due to resonance absorption or Raman scattering does not obey the condition λ_{mfp}/L_n , $L_T \ll 1$, except at very high background densities. Thus instead of describing the heat flux due to suprathermals by Eqs. (64) and (65), one is driven toward a description which explicitly recognizes the collisionless nature of the hot-electron component.

The issue of heat transport in laser-produced plasmas has gained prominence recently, because of a large body of experimental evidence at laser intensities $I_L \lambda_L^2 > 10^{14} \text{ W}\mu\text{m}^2/\text{cm}^2$ that the heat flux is well below the values in Eqs. (64) and (65) predicted by classical thermal conduction. The relevant experiments are reviewed by H. Ahlstrom in an accompanying report. Most of the experiments give data on the transport of thermal, rather than suprathermal, electrons. The experimental techniques include x-ray line spectroscopy [77], continuum x-ray spectra and target absorption [78], and the disassembly time of thin-foil targets [79]. Virtually all of these experiments imply thermal-electron heat transport that is significantly below the classical prediction of Eqs. (64) and (65), both for laser pulse lengths in the 100-ps range [77-79] and in the 1-ns range [80]. A few experiments have specifically addressed the transport of suprathermal electrons [81]. The latter data on K α emission are somewhat more difficult to interpret quantitatively, but in general the authors in Ref. [81] concluded that transport of hot electrons was not strongly inhibited. This result appears to be in conflict with those of Banner and Mead [77] and of Fabbro [79], who inferred that suprathermals were strongly inhibited. The apparent contradiction here clearly needs to be resolved.

Transport rates depend sensitively on knowledge of spatial gradients and time behavior. For example the typical evolution of the electron temperature obeys an equation such as

$$n_e \frac{\partial T_e}{\partial t} \approx - \frac{\partial}{\partial x} \kappa \frac{\partial T_e}{\partial x} \quad (67)$$

Thus in order to decide experimentally whether the value of κ predicted by Eqs. (64) and (65) is accurate, it is necessary to know the derivatives $\partial T_e / \partial t$ and $\partial T_e / \partial x$. However in most transport experiments [77-80] the data are integrated over either space, time, or both, so that the gradients are not yet measured directly. Hence the interpretation of the results is generally accomplished with the aid of numerical hydrodynamics computer models of the experiment, which provide a theoretical prediction of the values of $\partial T_e / \partial t$ and $\partial T_e / \partial x$ needed in Eq. (67).

Using such computer-hydrodynamics models, one can conclude from Refs. [77]-[80] that strong inhibition of thermal transport below the value predicted by Eqs. (64) and (65) seems to exist in both long and short-pulse irradiations, and in both high-Z and low-Z targets. In addition, the deduced transport seems to be inhibited both inward toward high-density regions, and laterally toward the exterior of the laser spot, as illustrated in Fig. 30. The quantitative value of the electron heat flux deduced from these experiments is roughly

$$q_e \lesssim - \left(\frac{\kappa}{20} \right) \nabla T_e \quad , \text{ or} \quad (68a)$$

$$q_e \lesssim - \left(\frac{1}{30} \right) (n_e T_e) v_{Te} \left| \frac{\nabla T_e}{\nabla T_e} \right| \quad (68b)$$

where κ is given in Eq. (65).

However, the reader should remember when using these numbers that they are only as reliable as the numerical-hydrodynamic models from which they were derived. When conduction is strongly inhibited, computer codes typically show very steep gradients of temperature and density. The numerical results then become sensitive to the finite size of the zones in space, and to the finite-difference algorithms used to solve the heat equation [1]. In addition, most of the experiments in Refs. [77]-[79] were done at high enough laser intensities that the heat flux carried by suprathermal electrons was probably the same order of magnitude as that carried by thermals. Because the numerical treatment of suprathermal electron transport in hydro codes is still somewhat incomplete in the collisionless regime, the influence of suprathermals on the inferred transport inhibition of the thermal population remains uncertain.

In the remainder of this section, we shall describe mechanisms which might be capable of reducing electron heat transport to the low values implied by experiments, Eq. (68). We shall treat in turn: large dc

FIG. 30. This figure points out the distinction between lateral heat flow in the plasma corona and inward heat flow from the critical surface toward high-density regions. The two processes may be influenced by rather different physical mechanisms.

magnetic fields, ion-acoustic turbulence, Weibel instabilities, and improved calculations of the "classical" values for thermal and suprathermal heat flux.

4.2 DC MAGNETIC FIELDS

Large-scale dc magnetic fields have been observed in laser-produced plasmas, with maximum amplitudes of a few times 10^6 gauss. The mechanisms responsible for producing such fields will be described in detail in Section 5. For the purposes of the present discussion, we shall confine ourselves to asking what kind of transport inhibition such megagauss-size magnetic fields might produce.

A well-ordered large-scale dc magnetic field slows electron transport across B when the electron Larmor radius $\rho_1 \equiv v_{te} m_e c / eB$ is less than the collisional mean free path λ_{mfp} . Under these conditions the step-size for random walk of the electrons across the field becomes ρ_1 instead of λ_{mfp} , as shown schematically in Fig. 31. The cross-field heat flux is reduced by the factor $1 + (\lambda_{mfp} / \rho_1)^2$ under these circumstances. Numerically, we recall from Eq. (12) that

$$\lambda_{mfp} / \rho_1 = n_e \tau_{ei} = 40 \left(\frac{T_e}{1 \text{ keV}} \right)^{3/2} \left(\frac{10^{21} \text{ cm}^{-3}}{Z n_e} \right) \left(\frac{B}{10^6 \text{ G}} \right) \left(\frac{10}{n_A} \right) \quad (69)$$

By looking for conditions where $\Omega_{ce} \tau_{ei} \gg 1$, this equation can be used to show when well-ordered megagauss-sized magnetic fields cause significant transport inhibition. We consider a moderate- Z target ($Z \sim 20$). For a laser with $\lambda_l = 1 \mu\text{m}$, Eq. (69) shows that a 1-MG field can inhibit transport of hot electrons with $T_e \gg 1 \text{ keV}$ even at densities of $2\text{--}3 n_c$. However it is harder to inhibit colder electrons, due to the $T_e^{3/2}$ dependence of λ_{ei} / ρ_1 . A thermal electron with $T_e \approx 1 \text{ keV}$ would only have $\Omega_{ce} \tau_{ei} \gg 1$ at densities above 10^{21} cm^{-3} if the magnetic field were greater than several megagauss. Conversely, a well-ordered 1-MG B field could inhibit electrons with $T_e \approx 1 \text{ keV}$ only at densities well below $n_c = 10^{21} \text{ cm}^{-3}$.

If we recall the distinction made in Fig. 30 between lateral transport and inward transport into dense material, we can conclude from the above discussion that a well-ordered 1-MG B field would be capable of inhibiting the lateral transport of both thermal and suprathermal electrons, at background densities below 10^{21} cm^{-3} . On the other hand transport into high-density material is more difficult to inhibit. Inhibition at

FIG. 31. Cross-field diffusion of an electron in a strong magnetic field. The effective step-size for the electron's random walk is the Larmor radius ρ_L .

densities of $2-3 \times 10^{21} \text{ cm}^{-3}$ is expected to be strong only for suprathermal electrons or for B fields closer to 10 MG.

4.3 ION-ACOUSTIC TURBULENCE

Large coherent fluctuations in the ion charge density are a second possible mechanism for transport inhibition. Short-wavelength ion fluctuations with spatial scales $\sim \lambda_{De}$ cause the electrons to scatter from coherent ion charge bunches instead of from individual ions. The resulting stronger scattering reduces the electron mean free path and the electron thermal conduction.

In the context of large electron heat fluxes, it has been proposed [82] that the electron heat flow itself could give rise to ion turbulence via the ion-acoustic instability. The physical basis for such a mechanism is illustrated in Fig. 32, which shows the ion and electron distribution functions in the presence of an electron heat flux q_e . The ions are taken to be considerably colder than the electrons, so their distribution is narrow. The electron distribution is distorted for $v_x < 0$ because heat is being carried in the $-x$ direction by electrons with speeds $|v_x| > v_{te}$. To avoid a net current flow, the electron distribution for $v_x > 0$ is skewed somewhat to the right, forming the "return current" of colder electrons.

FIG. 32. Electron and ion distribution functions in the presence of an electron heat flux q_e . When $u_s \gg c_s$ and $Z T_e \gg T_i$, the ion-acoustic instability grows rapidly.

Because of the return-current, the peak of $f_e(v)$ occurs at $v_s = +u_s$, rather than at $v_s = 0$. Between $v_s = 0$ and $v_s = +u_s$, the ions see a negative slope on the electron distribution function, and a Landau-type instability can result. Ion waves are unstable with a sizeable growth rate if $u_s > c_s$ and $c_s \gg v_{ti}$ (i.e., $Z T_e \gg T_i$). The wave vector $\mathbf{k}$ of the ion-acoustic waves is parallel to $-\mathbf{q}_e$.

The details of this ion-acoustic instability depend on the choice of initial distribution function $f_e(v)$. Forslund and Bickerton [82] studied the collisional case where the electron mean-free-path is small compared to the temperature gradient scale length, $\lambda_{mfp} \ll L_T$. These authors used a distribution function taken from collisional transport calculations [83], which gave an effective drift speed for the return-current equal to $u_s = 0.44 (\lambda_{mfp} / L_T) v_{te}$ for $Z = 1$. For $T_e \gg T_i$, the collisional distribution was predicted to be unstable when

$$u_s > c_s, \text{ or} \quad (70a)$$

$$\frac{\lambda_{mfp}}{L_T} > 2.3 \left(\frac{m_e}{M_i} \right)^{1/2}. \quad (70b)$$

In contrast, Manheimer [82] studied ion-acoustic instabilities in nearly collisionless systems. In this case there is no longer a unique self-consistent choice for $f_e(v)$, since in the absence of collisions an electron distribution can always carry more heat just by shifting a few electrons to very high velocities in the direction of $\mathbf{q}_e$. Manheimer chose a simple form of $f_e(v)$ which is nearly Maxwellian:

$$f_e(v) = \frac{n_0}{(2\pi v_{te})^{3/2}} \exp - \left\{ \frac{(v_x - u)^2 + v_y^2 + v_z^2}{2v_{te}^2} \left[1 + \frac{Q}{3n_e m_e v_{te}^3} \left(\left[\frac{v_x - u}{v_{te}} \right]^3 - 3 \left[\frac{v_x - u}{v_{te}} \right] \right) \right] \right\} \quad (71)$$

It is assumed that $q_e < (3/2)n_e m_e v_{te}^3$, and $Z = 1$. The first four moments of (71) are $n_e = n_0$, $\bar{v} = u\hat{x}$, $T_e = m_e v_{te}^2$, $q_e = Q\hat{x}$. Manheimer's distribution is unstable to ion-acoustic instability for

$$\frac{Q}{n_e m_e v_{te}^3} > c_s \left[1 + \left(\frac{M_i}{m_e} \right)^{1/2} \left(\frac{T_e}{T_i} \right)^{3/2} \exp - \left(\frac{T_e/2T_i + 3/2}{1 + k^2 \lambda_{De}^2} \right) \right] \quad (72)$$

for the case of equal electron and ion bulk velocities. When $T_e \gg T_i$ this is analogous to condition (70a) derived for the collisional distribution function.

From the point of view of transport inhibition in laser-heated plasmas, one must also know the effect of a *given* amount of ion turbulence on the plasma transport coefficients. In this latter area, no real theoretical consensus exists yet. Published estimates for the effective collision frequency ν_{ei} range from Manheimer's comparatively large value [82]

$$\nu_{eff} = \frac{1}{3} \left(\frac{\delta n_i}{n_i} \right)^2 \omega_{pe} \quad (73a)$$

to Lindman's very small value [84]

$$\nu_{eff} = 3 \times 10^{-5} \omega_{pe} \left[\frac{T_e}{T_i} \frac{u_s}{v_{te}} \left(\frac{m_e}{M_i} \right)^{1/4} \right] \quad (73b)$$

Here u_s is the drift speed shown in Fig. 31; $\delta n_i/n_i$ is the level of ion fluctuations, generally limited to $< 10\%$ by ion trapping [34,82]. Because of the large difference between (73a) and (73b), this is clearly an area where more computational and theoretical work is needed. The relevant computer simulations are reviewed in an article by Mead [85].

However, even if we use Manheimer's large value for ν_{eff} [Eq. (73a)], it seems that ion turbulence is not able to limit heat flow at densities above about 10^{21} cm^{-3} , the critical density for $\lambda_i \approx 1 \mu\text{m}$. The reasoning is as follows. Even with a very large turbulence level $\delta n_i/n_i$, it is numerically difficult to make ν_{ei} in (73a) larger than the classical collision frequency ν_{ei} . For example consider $\delta n_i/n_i$ as large as 0.1, and take $n_e = 10^{21} \text{ cm}^{-3}$, $T_e = 1 \text{ keV}$, $Z = 20$. Then from (73a) the ion turbulence produces an effective collision rate $\nu_{eff} \approx 6 \times 10^{12} \text{ s}^{-1}$, whereas the classical collision frequency is $\nu_{ei} \approx 8 \times 10^{12} \text{ s}^{-1}$. Under these conditions the turbulence has little or no effect, since $\nu_{ei} > \nu_{eff}$.

Furthermore, there is reason to think that $\delta n_i/n_i$ may not be as large as the 10% value used above. Langdon has emphasized the following dilemma. The growth rate for ion-acoustic instability is very slow near the threshold $(u_s)_{min} \approx c_s$. The growth-rate and angular spread of the ion waves do not become substantial until the drift speed u_s is an important fraction of the electron thermal speed, very roughly $u_s > v_{te}/4$. However, experiments showing inhibited transport indicate instead that $u_s \approx v_{te}/30 \approx 2c_s$, if we use Eq. (68b) with $A \approx 2Z$. This experimentally inferred value for the drift speed does not seem to be high enough to produce fast-growing ion-acoustic turbulence with a broad angular spectrum.

We thus emerge with the following (rather uncertain) picture of the role of ion turbulence in transport inhibition. The heat flux deduced from experiments seems to be large enough to exceed the thresholds for ion-acoustic instability, Eqs. (70) and (72). Thus one would predict theoretically that ion-acoustic turbulence should be produced, although the growth-rate and angular spread may not be large. Ion

turbulence has been observed [37] in recent CO₂ laser-plasma experiments, and it may also be produced by processes unrelated to electron heat flow, such as multi-species ion streaming instabilities [34,36].

The largest uncertainty arises when we try to estimate the effect of such ion turbulence on transport. Even the most generous estimates for the effective collision frequency are not big enough to compete with classical collisions at densities much above $10^{21} \mu\text{m}^{-3}$, unless Z is very small or the plasma above $n_c = 10^{21} \text{cm}^{-3}$ is very hot. If most of the laser light absorption occurs at densities $\geq 10^{21} \text{cm}^{-3}$, as would be the case for a) resonance absorption, $\lambda_1 = 1 \mu\text{m}$ or b) most absorption processes, $\lambda_1 \leq 0.5 \mu\text{m}$, we are therefore led to the conclusion that ion turbulence cannot explain the inhibition observed to exist between the absorption region and higher density regions.

However at densities below 10^{21}cm^{-3} and electron temperatures above a few keV, Eq. (73a) indicates that effective collisions due to ion turbulence might far exceed classical electron-ion collisions. These circumstances are most likely to occur in the *underdense* regions of plasmas produced by high-intensity CO₂ or Nd-glass lasers, which have critical densities $n_c < 10^{21} \text{cm}^{-3}$. One is thus led to the feedback scenario proposed by Manheimer and colleagues for this regime [35]: 1) ion turbulence in the underdense plasma absorbs laser energy, producing an inward heat flux toward n_c ; 2) the inward heat flux drives an ion-acoustic instability in the underdense plasma; 3) the ion-acoustic instability limits the heat flux, thus maintaining the necessary high plasma temperatures, and it also provides the ion fluctuations needed for the originally postulated good absorption. Ion turbulence that has been measured in the underdense CO₂ laser-produced plasmas of Ref. [37] may fit rather well into this sort of scenario.

4.4 WEIBEL INSTABILITIES.

Ion-acoustic turbulence is not the only process which can be driven unstable by electron heat flux. Anisotropy of the distribution function $f_e(v)$ due to the heat flux carried by the electrons can drive a Weibel instability, producing small-scale magnetic perturbations in the plasma [86]. These magnetic perturbations in turn are capable of reducing the electron heat flow by scattering the electrons.

The physics of a Weibel-type instability is best illustrated by the very simple example shown in Fig. 33. Consider a cold fluid, flowing in the z direction with speed $V_0 = V_0 \hat{z}$. Assume there are some initial perturbations in the x magnetic field, δB_x , which vary in the y direction:

$$\delta \mathbf{B} = \hat{x} \delta B_x \exp(ik_y y) \quad (74)$$

We shall show that because of the flow speed V_0 , these initial perturbations $\delta \mathbf{B}$ can grow exponentially.

The plasma electrons feel a $V_0 \times \delta \mathbf{B}$ force in the y direction, which pushes the electrons and creates a modulated electron density perturbation $\delta n \exp(ik_y y)$. Due to the flow velocity V_0 , the charge density results

FIG. 33. Simple example of a Weibel type instability. The plasma has a zero-order velocity $V_0 \hat{z}$, and an initial perturbation in the magnetic field $\hat{x} \delta B_x \exp(ik_y y)$. The $V_0 \times \delta \mathbf{B}$ force results in a perturbed current which enhances the original magnetic perturbation.

in a plasma current $\delta J \propto \hat{z} V_0 \delta n \exp(ik_y y)$. Finally, via Maxwell's equations the current δJ_z produces more of the original magnetic perturbations $\delta \mathbf{B}$, and thus the process grows exponentially.

The Weibel instability considered by Ramani and Laval must be treated by kinetic rather than fluid equations, but its basic physics is similar to the above simple example. The heat flux $\mathbf{q}_e$ replaces the cold-fluid flow velocity $\mathbf{V}_0$, and a perturbed distribution function δf_e replaces the perturbed electron density δn . The mechanism which produces magnetic perturbations is the same. Maximum growth occurs for perturbation wave-vectors $\mathbf{k}$ which are perpendicular to $\mathbf{q}_e$, just as the simple fluid example had $\mathbf{k} \perp \mathbf{V}_0$. The kinetic case was solved by Ramani and Laval in the limit [86]

$$\lambda_{\text{inst}} \ll \lambda_{\text{mfp}} \ll L_T \quad (75)$$

where the wavelength λ_{inst} of the unstable $\delta \mathbf{B}$ s is smaller than the electron mean free path, which is in turn smaller than the temperature gradient scale length.

The Weibel instability produces magnetic perturbations $\delta \mathbf{B}$ which vary in space on scales smaller than the Larmor radius of the electrons carrying the heat flux: $\rho_L = v_{te} m_e c / e \delta B > \lambda_{\text{inst}}$. Thus the heat-flux-carrying electrons are deflected by many small encounters with the magnetic perturbations $\delta \mathbf{B}$. Ramani and Laval's estimate for the consequent effective collision frequency is [86]

$$\nu_{\text{eff}} \approx \frac{v_{te}}{L_T} \left(\frac{L_T \omega_{pe}}{c} \right)^{1/3} \quad (76)$$

It is this effective collision frequency which determines the effective mean free path λ_{mfp} and scale length L_T to be used in Eq. (75).

If we compare the Weibel collision frequency (76) with the maximum predicted collision frequency due to ion-acoustic turbulence, Eq. (73a), we find that at $T_e = 2 \text{ keV}$, $n_e = 10^{21} \text{ cm}^{-3}$, $\delta n_i/n_i \approx 0.1$, the effective collision frequency due to Weibel is larger than the maximum estimate for ion-acoustic when $L_T < 10 \mu\text{m}$. The Weibel process has several theoretical advantages over ion-acoustic turbulence as a transport-inhibition mechanism: 1) Weibel persists even when $ZT_e \rightarrow T_i$, whereas ion-acoustic instability requires $ZT_e \gg T_i$; 2) The angular spectrum of ion-acoustic turbulence is rather narrow when the heat flux is small, and thus electrons with velocity nearly parallel to $\mathbf{q}_e$ are not strongly scattered by the ion waves. For Weibel, on the other hand, turbulent diffusion occurs for all electrons; 3) The actual value of ν_{eff} due to ion turbulence may be considerably smaller [84] than the maximum quoted in Eq. (73a). If this were the case, the collision rate due to Weibel in Eq. (76) would dominate that of ion-acoustic turbulence.

In summary, Weibel instabilities driven by electron heat flow seem theoretically promising as a way to inhibit transport. However, there are several theoretical questions which remain to be addressed. The saturation mechanism for the magnetic perturbations needs to be identified more definitively, and the relative importance of modes with $\mathbf{k}$ exactly perpendicular to the heat flux should be assessed. The influence of a background magnetic field on the growth of magnetic perturbations and on the subsequent scattering of electrons needs to be considered. Finally, the nonlinear response of background temperature and density gradients to the Weibel instability needs to be studied.

4.5 IMPROVED CALCULATIONS OF "CLASSICAL" HEAT FLUX

In this section we define "classical" transport as that due to electron-ion Coulomb collisions (perhaps in the presence of large-scale magnetic fields) *excluding* the effects of the various micro-instabilities and small-scale magnetic fields we have discussed above. We consider the possibility that if we are not calculating "classical" transport correctly, we may mistakenly conclude that an "anomalous" mechanism is needed to explain the low electron heat fluxes seen in experiments. To phrase this in a different way, Eq. (68a) indicates that heat flow is slower than the value $-\kappa \nabla T_e$, where the numerical value of κ is calculated by Spitzer [19,83] and Braginskii [10]. In the present section we ask whether the classical value of κ might actually be considerably lower than that derived by Spitzer and Braginskii, due to a variety of different physical effects. We

consider separately the heat flux expected for thermal "cold" and for suprathermal "hot" electrons, because the issues are somewhat different for the two cases.

4.5.1 Classical Heat Flux for Thermal Electrons

The classical calculation of the heat conductivity κ by Spitzer and Härm [83] utilizes an asymptotic expansion with small parameter λ_{mfp}/L_T , the mean free path divided by the temperature gradient scale length. For the expansion to converge it is necessary that λ_{mfp}/L_T be much less than unity.

Recently Gray and Kilkenny [37] investigated the values of λ_{mfp}/L_T for which Spitzer and Härm's calculation begins to fail. Their criterion for failure was that the electron distribution function $f_0(v)$ became negative for electrons with $|v|$ in the range $2.3v_{te} < |v| < 3v_{te}$, since it is these electrons which carry most of the heat flux. It was found that the Spitzer-Härm calculation can fail for values of λ_{mfp}/L_T as low as 0.015, or equivalently at values of $q_e/n_e m_e v_{te}^3$ as low as 0.03. Referring to Eq. (68), we see that this failure of the classical calculation occurs just when λ_{mfp}/L_T is nearly equal to the value 1/30 inferred experimentally. In this case it seems clear that an improved calculation of κ is needed, valid for values of the mean free path $0.015 L_T < \lambda_{\text{mfp}} < L_T$. Preliminary work in this area has been done by Bell and his colleagues [87], and indicates heat fluxes lower than those predicted by Spitzer and Härm.

A second physical process which might decrease the classical heat flux below the Spitzer-Härm value was suggested by Landon [18]. As described previously in Section 4.1, inverse Bremsstrahlung absorption can create a distinctly non-Maxwellian distribution function when $Z(e_L^2/n_e \omega_L^2) \gg v_{te}^2$. When this inequality is obeyed, Fig. 5 shows that the high-velocity tail of the distribution is lower than it would be for a Maxwellian. Since heat is carried in large degree by electrons with $v \approx 2-3v_{te}$, it is clear from Fig. 5 that the distribution produced by inverse Bremsstrahlung will carry less heat than an equivalent Maxwellian would. Although quantitative calculations of this effect in a temperature gradient have not yet been done, preliminary estimates [18] suggest that the heat flux may lie as much as a factor of two below the Spitzer-Härm value [83].

4.5.2 Classical Heat Flux for Suprathermal Electrons

The problem of finding a description of hot-electron transport that is both physically accurate and numerically tractable remains very difficult. As a consequence the role of suprathermal electrons in transport inhibition is still unclear, since refined calculations of hot-electron transport are difficult to come by even in the absence of micro-instabilities and magnetic fields. Recall that in most experiments where transport inhibition was deduced [77-79], a sizeable heat flux was present due to hot electrons. We are thus left to wonder whether the deduced thermal inhibition was real, or whether it appeared only to "cover up" inadequacies or numerical errors in the treatment of suprathermal-electron transport.

One main issue in hot-electron transport arises when the hot electrons are genuinely collisionless, so that $\lambda_{\text{mfp}} \gg L_n, L_T$. The diffusion approximation, which is probably the most numerically efficient hot-electron description [88], is no longer valid in this regime because there is nothing to keep the electron distribution nearly isotropic. In particular, a corona dominated by collisionless hot electrons can at times interact only weakly with the dense collisional core of the target [89,90]. Consider the geometry shown in Fig. 34, where a large hot-electron corona of radius R surrounds a spherical core of dense material [90]. When the coronal electrons have a mean free path λ_{mfp} which is large compared to the radius R , they can bounce around in the corona many times before hitting the core. Under these conditions the heat flux into the dense material is [90]

$$|q_e| \cong \left(\frac{R}{\lambda_{\text{mfp}}} \right) n_e m_e v_{te}^3, \quad \frac{R}{\lambda_{\text{mfp}}} \ll 1. \quad (77)$$

(Recall that in the classical expression for the heat flux, Eq. (63), it is the *inverse* of R/λ_{mfp} which appears.) If one were using a phenomenological "flux-limited" expression for the heat flux, such as

$$|q_e| \cong f n_e m_e v_{te}^3, \quad (78)$$

one would conclude that the "flux-limit" f was $R/\lambda_{\text{mfp}} \ll 1$ in this situation. This is suggestive of the form (68b) derived from experiments, where $f \approx 1/30$.

FIG. 34. Geometry for corona-core decoupling [89,90]. If the solid angle of the dense-plasma core is small as seen by a hot coronal electron, the electron can bounce many times in the corona before hitting the core and losing its energy there.

A second effect related to the numerical treatment of hot electrons concerns the outer edge of the corona. In low-density regions, theory predicts [4] that the hot-electron density can exceed that of the cold electrons. This is because the hot electrons expand faster than the colds, and because they can penetrate farther into the electrostatic potential sheath at the edge of the plasma before being turned around. Such an effect is particularly important for CO_2 lasers and high-intensity Nd-glass lasers. It is generally not well described in the usual diffusive hot-electron schemes, because a Poisson equation is needed for the potential. Recently some numerical treatments have appeared which can specifically include this hot-electron dominated region in large hydrodynamics codes [91,92].

Generally under laser-fusion conditions, the time scales characterizing electron heat transport are short compared to overall hydrodynamic time scales. Thus one can afford to do much more sophisticated transport studies by keeping the hydrodynamic variables $n(\mathbf{x})$ and $v(\mathbf{x})$ fixed, and then solving in detail for the time and space evolution of heat flow due to hot and cold electrons. Recently Mason [93,94] has used such a technique to develop transport models which are rather more elaborate than had previously been possible in fully hydrodynamic computer codes. Mason's work on the self-consistent role of hot electrons in magnetic field generation [93] showed the effect of a large-scale magnetic field around the edge of the laser spot on lateral transport of thermals and suprathermals. This work will be discussed in the B-field discussion, Section 5.

Mason's second study, Ref. [94], concerned an issue first posed several years ago. Valeo and Bernstein [95] suggested that hot-electron penetration into dense matter could be inhibited by the need to draw a neutralizing return current through the cold, collisional background plasma. Roughly speaking, the zero-current condition between hot and cold electrons is

$$n_h v_h = n_e v_e \quad (79)$$

so that the electric field due the cold-electron collision rate ν_c is given by

$$E = - \left(\frac{m_e}{e} \right) \nu_c v_e = \left(\frac{m_e}{e} \right) \nu_c \left(\frac{n_h}{n_e} \right) v_h \quad (80)$$

Kruer [96] has pointed out that expression (80) acts as if the *hot* electrons experience an effective collision rate

$$\nu_{\text{eff}} = \left(\frac{n_h}{n_e} \right) \nu_c \quad (81)$$

which can be substantially higher than their classical collision frequency. Thus hot electron transport appears “inhibited.”

Mason reexamined this process using a transport model with Monte Carlo hot electrons and fluid cold electrons [94]. He found the predicted inhibition effect to be strong when the cold electron temperature was held fixed at a value $T_e < 100$ eV. However when realistic descriptions of Joule heating and cold-electron thermal conduction were included to heat the thermals, it was found that the cold electron temperature T_e rose rapidly. Then because $\nu_c \propto T_e^{-3/2}$, the hot-electron transport inhibition was extinguished after times of the order of picoseconds. This effect has been also discussed by K. Brueckner.

With the same Monte Carlo hybrid code, Mason was able to isolate what he terms a “convective” mechanism for transport inhibition [94], due to cold-electron convection and thermo-electric terms in Braginskii’s [10] heat equation. Because of the quasineutrality condition (79), large cold-electron convective heat fluxes $n_e v_e T_e$ arise when there are large hot-electron fluxes $n_h v_h$. The effect is enhanced when there are “density troughs” in the overdense plasma (the latter have been predicted theoretically [97] when the plasma blowoff is already supersonic at densities above n_c). *Strong inhibition of cold-electron transport* was seen computationally in the presence of a “density trough,” whereas without such a trough the convective and thermoelectric terms were not large enough to create significant inhibition.

To summarize our discussion of “classical” transport, it seems clear that there are several ways in which both thermal and suprathermal electron transport have previously been overestimated, even in the absence of micro-instabilities and small-scale magnetic fields. No single one of these effects presently seems large enough to explain the observed factor of 20–30 decrease in heat transport below the conventional values. Yet I think it is entirely possible that the combination of several of these mechanisms may in the end eliminate the “discrepancy” between experimental and theoretical values for the electron heat flux. I would strongly recommend this type of issue as a fruitful area for future research.

5. PRODUCTION OF LARGE DC MAGNETIC FIELDS

5.1 EXPERIMENTAL EVIDENCE

DC magnetic fields have been observed in laser-produced plasmas. For many years these fields were measured far from the critical surface, by magnetic probes. More recently, Faraday rotation of a high-frequency electromagnetic probe beam has been used [98,99]. For some target configurations, this technique has allowed fields of several megagauss (several $\times 10^6$ G) to be inferred to exist in the plasma corona.

For a linearly polarized probe beam of frequency ω , the position angle of the polarization vector is rotated by an angle

$$\Delta\theta \approx \frac{\omega}{2c} \int_0^L \frac{\omega_p^2}{\omega^2} \frac{c\mathbf{B}}{m_p c \omega} \cdot d\mathbf{r} = \int n_e \mathbf{B} \cdot d\mathbf{r} \quad (82)$$

after traversing a path length L through the plasma. Thus by measuring the Faraday rotation angle $\Delta\theta$, one learns the average of $n_e \mathbf{B}$ along the path of the probe beam. If one can simultaneously use interferometry to measure the average density n_e along the same path, a mean value of $\mathbf{B}$ can be obtained from

$$\langle \mathbf{B} \rangle = \frac{\int n_e \mathbf{B} \cdot d\mathbf{r}}{\int n_e d\mathbf{r}}$$

To obtain the *maximum* value of $\mathbf{B}$ along the ray path, one must accurately model the density profile and the refraction of the probe beam.

The geometry for a typical Faraday rotation experiment is shown in Fig. 35. To date the large-scale $\mathbf{B}$ fields which have been measured have been toroidal in nature, as indicated in the figure. This would produce opposite senses of Faraday rotation on the top and bottom rays of the probe beam, and no rotation of the middle rays. Clearly any small-scale inhomogeneities or turbulent structure in the magnetic field would not be measurable by this technique, since they would average out over the line of sight.

For laser pulses shorter than 100 ps, experiments with Nd-glass lasers indicate [98] the presence of a toroidal magnetic field which peaks in the underdense plasma ($n_e \approx 0.2 n_c$). However such fields are observed

FIG. 35. Geometry for Faraday-rotation measurement of the toroidal magnetic field on a disc target.

to exceed 100–200 kG only for targets larger than about 100 μm in diameter, and only when the target diameter is at least several times larger than the focal spot diameter. These conditions suggest to Raven et al. [98] that the electrical capacity of the target may play an important role in short-pulse experiments. Small isolated targets have low capacitance. Thus the escaping hot electrons from resonance absorption can rapidly charge the target to a high potential, with the resulting electric field inhibiting the further growth of current and B field. This hot-electron mechanism for magnetic field generation will be discussed in more detail in Section 5.2.2 below.

For laser pulse lengths longer than 1 ns, experiments [99] show a larger-scale toroidal magnetic field surrounding the laser spot. Preliminary results of Willi and Rumshy [99] indicate megagauss-sized fields under a wide variety of experimental conditions. One is tempted to attribute these toroidal fields to the $\nabla n \times \nabla T$ mechanism, discussed in Section 5.2.1 below. The long-pulse experiments also show filament-like structures [99] which modulate the Faraday probe beam, suggesting smaller-scale solenoid-shaped magnetic fields surrounding the cylindrical filaments. These will be discussed further in Section 5.3.1 below.

From the theoretical point of view, an embarrassingly wide variety of mechanisms seems capable of generating large dc magnetic fields in laser-produced plasmas. The list to date is as follows:

- 1) Mechanisms resulting in large-scale B fields
 1. $\nabla n_e \times \nabla T_e$ mechanism
 2. Hot-electron electron from focal spot
 3. Anisotropic electron pressure
- 2) Mechanisms producing small-scale B fields
 1. $\nabla n_e \times \nabla T_e$ fields near filaments, composition jumps, shocks
 2. B fields due to resonance absorption
 3. Thermal instabilities
 4. Weibel instabilities

We shall discuss each of these mechanisms in turn, although in varying amounts of detail. In the final part of this section we shall describe how electron transport might be affected if several of these types of magnetic fields were present in the plasma at the same time.

We begin with mechanisms capable of generating large-scale B fields, which extend over the scale of the laser spot or the target size.

5.2 MECHANISMS RESULTING IN LARGE-SCALE B FIELDS

5.2.1 The $\nabla n_e \times \nabla T_e$ Mechanism

This process was originally proposed by Biermann in 1950, in the context of rotating magnetized stars [100]. It was applied to laser-produced plasmas by Stamper and colleagues in 1971 [100]. Further work on this mechanism is described in Ref. [101]. The simplest explanation of the effect is as follows. In an expanding plasma, the electron pressure force gives rise to a charge-separation electric field

$$en_e \mathbf{E} = -\nabla p_e. \quad (83)$$

By Maxwell's equations, there is a source of B when this electric field is solenoidal:

$$\frac{\partial \mathbf{B}}{\partial t} = -c \nabla \times \mathbf{E} = -\frac{c}{e} \nabla \ell n n_e \times \nabla T_e. \quad (84)$$

Thus this mechanism generates a dc magnetic field when there is an angle between the local density and temperature gradients in the plasma.

For a target irradiated with one finite-sized laser beam, this is most likely to occur in the outer regions of the laser spot, as shown in Fig. 36. The density gradient points into the solid-density target, whereas near the edge of the laser spot the temperature gradient points radially inward. The result is a toroidal $\nabla n_e \times \nabla T_e$ magnetic field around the outer parts of the laser spot. Since the intensity distribution of the laser light

FIG. 36. For a finite laser focal spot, the prevailing temperature and density gradients produce a toroidal de magnetic field, by the $\nabla n_e \times \nabla T_e$ mechanism.

is usually rather flat near the middle of the focal spot, radial temperature gradients are thought to be small there. Thus one expects a region of low magnetic field at the center of the focal region.

To estimate the steady-state size of the $\nabla n_e \times \nabla T_e$ magnetic field, we can balance the B field growth rate (84) with convection due to outward fluid motion at velocity V :

$$\frac{c}{e} \nabla t n n_e \times \nabla T_e \approx \nabla \times (V \times B) \quad (85)$$

Apart from angular factors, this gives a steady-state magnetic field

$$B \approx \left(\frac{c}{c_s} \right) \left(\frac{T_e}{e L_T} \right) \quad \text{or} \quad (86a)$$

$$\left(\frac{B}{10^6 \text{ gauss}} \right) = 30 \left(\frac{T_e}{1 \text{ keV}} \right)^{1/2} \left(\frac{1 \mu\text{m}}{L_T} \right) \left(\frac{A}{Z} \right)^{1/2} \quad (86b)$$

Hence a 1-keV plasma with a temperature scale length of $20 \mu\text{m}$ would produce a magnetic field $B \approx 1.5 \text{ MG}$, at places where the density and temperature gradients were at large angles relative to each other.

In principle, resistive diffusion might also saturate the growth of B. However in practice, convection is generally larger numerically than resistive diffusion for all but the shortest gradient lengths:

$$\frac{\left(\frac{c^2}{4\pi\sigma} \right) \nabla \times B}{V \times B} = \frac{c}{c_s} \left(\frac{\nu_{ei}}{\omega_p} \right) \frac{c}{\omega_p L_n} = \frac{Z}{10} \left(\frac{1 \text{ keV}}{T_e} \right)^2 \left(\frac{1 \mu\text{m}}{L_n} \right) \quad (87)$$

$\nabla n_e \times \nabla T_e$ magnetic fields are a natural candidate for inclusion in numerical hydrodynamics codes [88,102,103], because of their geometrical simplicity and potential for explaining transport inhibition. However even in the absence of a separate hot-electron population, calculation of B fields can be quite a difficult undertaking from a numerical point of view. Colombant and Winsor [104] pointed out that the full equation for the evolution of the magnetic field [10] contains four other terms in addition to the $\nabla n_e \times \nabla T_e$

source term: convection, resistive diffusion, magnetic curvature and magnetic pressure, and the so-called "thermal force" [10]

$$- \frac{e}{c} \nabla \times [\beta \nabla T_e - \alpha \hat{b} \times \nabla T_e] \quad .$$

where α and β depend on (λ_{mp}, p_1) . They found that once the magnetic field begins to grow, all five terms in the magnetic-field equation are of comparable size, including the thermal force term. This is unfortunate, because it is difficult numerically to find an efficient differencing scheme for the full magnetic-field equation. (Various authors had attempted to solve it by neglecting one or more of the terms.)

Accurate numerical treatment of the $\nabla n_e \times \nabla T_e$ magnetic field becomes even more difficult when a hot-electron population is included. There are two aspects to this issue. First, consider the effect of the B field on the hot-electron fluxes. A collisionless plasma responds to an applied force F by drifting in a direction perpendicular to both B and F , at a speed proportional to $F \times B/B^2$. Hence in multi-group treatments of hot electrons, the cross-field drifts have the effect of mixing together the different spatial components of the hot-electron fluxes. The result is to make it very costly to include such drifts in a computer calculation when there are many electron energy groups to be followed. Mason [93] has been able to solve an idealized version of this problem self-consistently by treating the hot electrons as a separate Maxwellian fluid. The full multi-energy-group problem remains unsolved.

The second issue concerns the effect of hot electrons on the size of the magnetic field [105,106]. It is clear that hot-electron currents can play an important role in determining the overall structure of B . Even in the most elementary example, the $\nabla n_e \times \nabla T_e$ source term includes both cold- and hot-electron contributions:

$$\nabla n_e \times \nabla T_e = \nabla(n_c + n_h) \times \nabla(T_c + T_h)$$

Yet to date such effects are not generally included in numerical hydrodynamics codes, with the exception of Mason's transport work [93].

As a result of these practical issues, the inclusion of $\nabla n_e \times \nabla T_e$ magnetic fields in numerical hydrodynamics codes remains a difficult problem in the presence of suprathermal electrons. The hydro codes with the most complete B -field models do not in general include hot-electron transport. On the other hand, Lagrangian numerical hydrodynamics codes which include suprathermal electrons and which are used routinely to model experiments [48,103,107] do not yet have fully self-consistent treatments of the magnetic field interaction with hot electrons. Finally, the one transport code to date which does treat B fields and hot electrons self-consistently [93] does not yet do hydrodynamics. I hope the reader will conclude from the above discussion that the modelling of $\nabla n_e \times \nabla T_e$ magnetic fields coupled with hot electrons in numerical hydro codes is a challenging and important research area.

5.2.2 Hot-Electron Ejection from the Focal Spot

Hot electrons produced by resonance absorption are accelerated preferentially in an outward direction, i.e., down the density gradient. Several years ago it was pointed out [108] that if these hot electrons caused a return current to form in the colder regions outside the focal spot, a toroidal magnetic field could result. Dawson called this the "fountain effect" [108], because of the geometry which is illustrated in Fig. 37. The sense of the magnetic field is the same as that in the $\nabla n_e \times \nabla T_e$ mechanism described above. This concept has been discussed recently by Kolodner and Yablonovitch [109].

Present understanding of the fountain effect mechanism remains only qualitative, because the B field is sensitive to details of the return-current spatial structure. Clearly the magnetic field will be small if the hot electrons are able to draw a return current locally equal and opposite to J_{hot} while they are still in the focal spot. This might be the case near the center of the spot, for example. The "fountain effect" is strongest near the spot edge, and thus would be expected to have the largest global effect for experiments with small focal spots. It could dominate in short-pulse, high-intensity experiments [98,109].

Drouet [110] has used wire probes to measure return-current patterns similar to those predicted by Fig. 37.

FIG. 37. Configuration producing a "fountain effect" magnetic field due to hot-electron ejection from a finite focal spot [108,109]. The details of such a field are sensitive to the exact return-current path and the target capacitance.

5.2.3 Anisotropic Electron Pressure

Even when the electron density and temperature gradients are parallel, large-scale magnetic fields can be produced in the corona when the electron pressure tensor is not isotropic. This effect has been discussed in detail by Mora and Pellat [76].

The physics behind it is analogous to that of the $\nabla n_e \times \nabla T_e$ mechanism: In a steady-state expansion, the electron pressure creates an electric field

$$\mathbf{E} = -\frac{1}{en_e} \nabla \cdot \mathbf{P}_e, \quad (88)$$

where $\mathbf{P}_e$ is the full electron pressure tensor. The growth of magnetic field is then

$$\frac{\partial \mathbf{B}}{\partial t} = -c \nabla \times \mathbf{E} = \frac{c}{e} \nabla \times \frac{1}{n_e} \nabla \cdot \mathbf{P}_e. \quad (89)$$

Thus even if ∇n_e is parallel to ∇T_e , a B field can be produced when the pressure-tensor is not diagonal: $\mathbf{P}_e \neq P\mathbf{I}$. Mora and Pellat's kinetic-theory treatment of this mechanism [76] shows that B fields in the megagauss range can be produced in the underdense plasma, and that p-polarization of the laser beam is preferred.

The three processes discussed above are capable of creating magnetic fields with large-scale structure, on the scale of the laser spot itself. The total large-scale shape is predicted to be toroidal, due to the $\nabla n_e \times \nabla T_e$ and "fountain effect" mechanisms. Because of the anisotropic-pressure mechanism, this toroidal B field will be locally stronger in regions of the incident laser beam which correspond to p-polarization. The Faraday-rotation experiments of Ref. [98] measured large-scale fields which were toroidal in shape, as predicted. However, our ability to model these fields quantitatively is not refined enough to learn the exact mixture of the large-scale mechanisms responsible for the observed field.

5.3 MECHANISMS PRODUCING SMALL-SCALE B FIELDS

In addition to the three processes just discussed, there have been quite a few mechanisms proposed which would make Bs with small spatial scales, ranging from a few vacuum laser wavelengths to an electron mean free path. Because they lack large-scale coherence, such fields are difficult to observe by present

Faraday-rotation techniques which integrate along the probe's line of sight. However, we shall show in Section 5.4 that if these small-scale magnetic perturbations do exist, they can have a significant effect on the expected electron transport coefficients.

5.3.1 $\nabla n_e \times \nabla T_e$ Fields near Filaments, Composition Jumps, and Shocks

It is clear that if the laser-produced plasma contains density or temperature gradients with small spatial scales, the $\nabla n_e \times \nabla T_e$ mechanism described in Eqs. (84)–(86) can produce small-scale $\mathbf{B}$ fields as well as large-scale fields. We showed an example in Fig. 25 of a magnetic field surrounding a filament. In this case the local density gradient length is the width of the filament, which might be as small as a few vacuum laser wavelengths. Thus the resulting $\mathbf{B}$ -field would have a correspondingly small spatial scale. Such fields may have been observed in Ref. [99].

In a similar vein, Tidman [111] has suggested that the strong local gradients near composition jumps or shock waves could produce $\nabla n_e \times \nabla T_e$ magnetic fields. In the case of very sharp composition jumps, for example, Tidman estimated that the spatial width of the $\mathbf{B}$ field would be determined by sideways resistive diffusion to be $\Delta x \approx c \Delta t (\pi \omega)^{1/2} \approx$ a few microns, for a 100-ps high-power laser irradiation, $T_e \approx 1$ keV. Raven et al. [98] have measured the large field produced by a composition jump.

5.3.2 Magnetic Fields due to Resonance Absorption

Stamper and Tidman [112,113] proposed many years ago that radiation pressure due to laser light absorption could produce a dc magnetic field. Under these circumstances one can imagine replacing the general pressure tensor $\mathbf{P}_e$ discussed in Eqs. (88)–(89) by the electromagnetic stress tensor $\mathbf{P}_R$. In the presence of absorption, $\nabla \times (\nabla \cdot \mathbf{P}_R)$ does not vanish for light of the appropriate polarization, and a dc magnetic field can result.

Resonance absorption has proven to be an interesting example of this mechanism, because megagauss-sized dc magnetic fields have been seen to accompany resonance absorption in computer simulations [114,115]. Theory for the saturation of this process [76,115] shows megagauss-sized fields due to resonance absorption concentrated in a very thin region just at the critical-surface layer. Because the region containing the large magnetic field is less than a vacuum wavelength thick, $\mathbf{B}$ fields due to resonance absorption are not good candidates for explaining the inhibition of heat flow toward high densities [116]. Rather, they probably add to the general level of small-scale perturbations superposed on the overall toroidal magnetic field structure.

5.3.3 Fields due to Thermal Instabilities

The magnetic field dependence of the electron thermal conductivity has been shown to cause a field-generating thermal instability [117]. The physical mechanism for instability is that a perturbation in the plasma temperature δT , which varies across the zero-order density gradient, produces a magnetic field $\delta \mathbf{B}$ via a $\nabla n_e \times \nabla \delta T$ source term. The magnetic field dependence of the electron thermal conductivity then changes the heat flow in such a way as to enhance the original temperature perturbation δT , giving a feedback mechanism for instability. This effect requires the instability wavelength to be larger than the collisional mean free path, so it is most likely to be important in moderate- Z or high- Z plasmas.

The wavelength of the magnetic perturbation giving maximum growth in Tidman and Shanny's [117] theory is roughly equal to

$$\lambda_{\max} \sim \left(\frac{c}{\omega_p \lambda_{\text{mfp}}} \right)^{1/2} \left(L_n L_T \right)^{1/2} \quad (90)$$

when $Z \gg 1$. Here λ_{mfp} is the mean free path given in Eq. (11). Thus near the critical density for a Nd-glass laser, one obtains typical scale sizes for the magnetic perturbations equal to a fraction of the mean gradient length $(L_n L_T)^{1/2}$. As a result, thermal-instability magnetic fields have small spatial scales relative to the typical gradient length or spot size. They presumably maximize at densities slightly below critical, where the density and temperature gradients are in the same direction. The nonlinear behavior of this instability has not yet been quantitatively addressed.

5.3.4 Magnetic Fields due to Weibel Instabilities

Electromagnetic instabilities driven by inward heat-flow [86] can also produce small-scale magnetic perturbations in the plasma. Except for the singular limit where $\mathbf{k}$ is exactly perpendicular to the heat flux, these are not strictly speaking dc magnetic fields. Rather, they oscillate in time. However such fields can appear quasistatic to a fast-moving electron, if the electron's transit-time across $\delta\mathbf{B}$ is shorter than the oscillation timescale.

We can estimate the magnitude of the magnetic perturbations from the nonlinear theory of Ramani and Laval [86]. One can deduce the saturation amplitude from the effective collision frequency

$$\nu_{\text{eff}} \approx \sum_{\mathbf{k}} \left(\frac{e\delta B_{\mathbf{k}}}{m_e c} \right)^2 \frac{1}{k' v_{te}} = \frac{v_{te}}{L_n} \left(\frac{L_n \omega_p}{c} \right)^{1/3} \quad (91)$$

The most unstable mode has wavenumber

$$k \approx \frac{\omega_p}{c} \left(\frac{c}{L_n \omega_p} \right)^{1/3} \quad (92)$$

These two equations can be combined to yield

$$\frac{(\delta B)_{\text{max}}}{10^6 \text{ gauss}} \approx 2 \left(\frac{n_e}{10^{21} \text{ cm}^{-3}} \right)^{1/4} \left(\frac{T_e}{1 \text{ keV}} \frac{1 \mu\text{m}}{L_n} \right)^{1/2} \quad (93)$$

Once again we find a mechanism capable of producing megagauss magnetic fields. The spatial scale of these fields is quite small, since the analysis of Ramani and Laval requires instability wavelengths

$$\lambda_{\text{inst}} < \lambda_{\text{mfp}} < L_n$$

where λ_{mfp} is the mean free path.

5.4 HEAT TRANSPORT IN REALISTIC B-FIELD GEOMETRIES

Most of the mechanisms for magnetic field generation described above are theoretically capable of producing field magnitudes of at least a few hundred kilogauss. Several mechanisms are predicted to make fields as high as a few megagauss, for typical parameters of a laser fusion target. The overall field geometry (Section 5.2) is predicted to be dominantly toroidal on the large spatial scale, extending over distances typical of the laser spot-size. Superposed on this large-scale field, one expects to find small-scale fluctuations due to the processes described in Section 5.3. Typical spatial scales for the small-scale fields range from a fraction of the laser wavelength up to a fraction of the density or temperature gradient scale length. Thus the total magnetic field might resemble that shown [69] in Fig. 38.

(A parenthetical note of caution is appropriate at this point. The work to date on magnetic field generation has considered only one physical mechanism at a time. A potentially fruitful area for research concerns the interaction and interference of several different field-generation mechanisms, if they all occur simultaneously.)

Since laser-produced plasmas typically have very high values of $\beta = 8\pi n k T / B^2$, the magnetic field cannot directly affect the hydrodynamics, particularly since an ion Larmor radius is large compared with the scale length. The principal effect of the magnetic field is on the electron thermal conduction. Here we shall obtain approximate expressions for the cross-field diffusion coefficient of test electrons in realistic B-field geometries. It is important to realize that this diffusion of test electrons does not correspond to a

FIG. 38. Schematic representation of total dc magnetic field produced by laser-plasma interaction, in the limit that the field lines are not very stochastic. Jokipii [119] called such a configuration the "wet spaghetti" model for a turbulent magnetic field.

hydrodynamic motion of electrons. If the electrons begin to drift away from the ions, it is only too easy for electric fields to be generated which bring them back. However, the diffusion of test electrons does correspond to a thermal conductivity. The reason is that when electric fields cancel particle fluxes induced by diffusion, they do not cancel energy fluxes. Thus, the test particle diffusion coefficients discussed in this section are a reasonable approximation to the electron thermal conduction.

Cross-field electron energy transport is well understood in situations when the field lines are nicely ordered and nearly straight. Figure 31 illustrates the particle diffusion process when the electron cyclotron frequency is large compared to the collision frequency, $\Omega_{ce} \gg \nu_{ei}$. The step-size for cross-field diffusion is the electron Larmor radius ρ_L , and the characteristic collision time is ν_{ei}^{-1} , so the "classical" particle diffusion coefficient across the field is

$$D_{ei} \approx \rho_L^2 \nu_{ei} = v_{te}^2 \nu_{ei} / \Omega_{ce}^2 \quad (94)$$

This is $(\Omega_{ce} / \nu_{ei})^2$ times smaller than the diffusion coefficient for $B = 0$.

However, we have argued above that the field lines in laser-fusion targets are not likely to be particularly well-ordered and straight. In this case one might prefer to describe the total field as a large-scale ordered field $\mathbf{B}_0(\mathbf{x})$, with some smaller-scale fluctuations $\delta\mathbf{B}(\mathbf{x})$ superposed:

$$\mathbf{B} = \mathbf{B}_0(\mathbf{x}) + \delta\mathbf{B}(\mathbf{x}) \quad (95)$$

In this case the presence of small-scale fluctuations $\delta\mathbf{B}$ can act to enhance the cross-field diffusion over its classical value (94).

This issue of heat transport in a fluctuating dc magnetic field has been studied rather extensively in other contexts. Two different regimes present themselves, depending on the degree of disorder of the total field. If the fluctuations $\delta\mathbf{B}$ are large enough and if they have appropriate scale-sizes, the total field can be strongly stochastic: two field lines which are close to each other at one point in space can wander away from each other randomly, and eventually fill a sizeable fraction of the available volume. In this limit, electrons moving along the stochastic lines can lead to *transverse* heat flow, due to their motion *parallel* to the wandering magnetic field. Figure 39 is a crude illustration of such a stochastic situation.

FIG. 39. Stochastic magnetic field lines. When the small-scale magnetic perturbations are large enough and of appropriate scale sizes, the overall field line structure can become stochastic: two field lines adjacent to each other at one point in space can wander arbitrarily far apart at other spatial locations. Electron transport in this type of field structure will be qualitatively different from that in the more regular structure shown in Fig. 38.

Several authors have investigated electron transport in the stochastic magnetic fields characteristic of tokamak geometry [118]. The application of tokamak-related work to the laser-plasma situation is not immediate, because in the latter case there is little magnetic shear and the plasma is high-beta ($B^2/8\pi \ll nT$). However, it is intuitively evident that field-line stochasticity must play a role in laser-produced plasmas as well as in tokamaks, if the magnetic field fluctuations around the toroidal mean field are large enough and of appropriate spatial scales. This is a promising area for further research.

In the opposite limit the magnetic-field fluctuations are sufficiently small and well-ordered that the total field retains a roughly coherent shape, as suggested in Fig. 38. Particle diffusion in such fields has been studied extensively in the context of cosmic-ray propagation [119].

Applications to laser-produced magnetic fields were discussed in Ref. [69]. The results are shown schematically in Fig. 40, and may be summarized as follows. The general effect of small-scale magnetic field fluctuations is to make cross-field heat flow more rapid than the classical rate for a nonfluctuating $\mathbf{B}$ field, $D_{cl} \approx \mu_1^2 \nu_{ei}$. Consider a total magnetic field given by

$$\mathbf{B} = \mathbf{B}_0 + \sum_{\mathbf{k} \neq 0} \delta \mathbf{B}_{\mathbf{k}} \quad , \quad (96)$$

and assume for the sake of simplicity that $\sum_{\mathbf{k}} |\delta \mathbf{B}_{\mathbf{k}}| \ll |\mathbf{B}_0|$. For very small values of the electron-ion collision frequency, $\nu_{ei} \ll kv_{te}(\delta \mathbf{B} \cdot \mathbf{B}_0)^{1/2}$, electrons can bounce many times in small magnetic mirrors produced by the $\delta \mathbf{B}$'s. Cross-field drift motion during each bounce then leads to a "neoclassical" diffusion rate [69]

$$D_{nc} \approx D_{cl} \left[\sqrt{\frac{\delta B}{B_0}} \frac{\pi^2}{2} \frac{k_{\perp}^2}{k_{\parallel}^2} \right] \quad , \quad (97)$$

where $k_{\perp}$ and $k_{\parallel}$ are the components of $\mathbf{k}$ perpendicular and parallel to $\mathbf{B}_0$.

FIG. 40. Cross-field electron diffusion coefficient [69] D for a turbulent magnetic field such as that of Fig. 38, versus collision frequency ν_{ei} . Parameters are $\delta B/B_0 = 0.36$, and in the neoclassical regime $\lambda = \lambda_L$, and $\mu_L/\lambda = 0.04$.

For slightly larger values of the electron-ion collision frequency, $k v_{te} (\delta B/B_0)^{3/2} < \nu_{ei} < \Omega_{ce} (\delta B/B_0)^2$, two quasilinear processes dominate the cross-field heat flow. If the fluctuation scale-size parallel to $\mathbf{B}_0$ is comparable to an electron Larmor radius, $\lambda_z \approx 2\pi\mu_L$, electrons moving along the field lines are deflected by the $\mathbf{v} \times \delta \mathbf{B}$ force once every gyro orbit. This leads to a cross-field diffusion rate due to resonant scattering [69]

$$D_{rs} \approx \left(\frac{v_{te}^2}{\Omega_{ce}} \frac{\sum_k \delta B_k^2}{B_0^2} \right) \quad (98)$$

If δB s are *not* present with the appropriately resonant parallel wavelengths, diffusion is instead determined by the transverse "wandering" of the field lines described by Jokipii and Parker [119]:

$$D_w \approx v_{te} L_c \left(\frac{\sum_k \delta B_k^2}{B_0^2} \right) \quad (99)$$

Here L_c is the correlation length of the fluctuations.

Finally, in the more collisional limit $\Omega_{ce} (\delta B/B_0)^2 < \nu_{ei} < \Omega_{ce}$, cross-field transport is the classical value for a magnetized plasma $D_{cl} = \rho_L^2 \nu_{ei}$. These three regimes of collisionality are depicted in Fig. 40 for plasma parameters listed in the figure caption. (The resonant scattering expression was used for the "quasilinear" regime.)

Numerical evaluation of the cross-field diffusion coefficients given in Eqs. (97)–(99), under plausible laser-plasma conditions [69], shows that these physical mechanisms produce transport that is still considerably slower than the classical value for a magnetic field-free plasma $D_{B=0} = v_{te}^2 \nu_{ei}$. Thus a fluctuating but non-stochastic laser-produced $\mathbf{B}$ field seems capable of explaining some part of the transport inhibition inferred from experiments, at least at plasma densities less than critical. Stochasticity would increase the effective heat flow across the field, but the magnitude of this effect has not yet been studied.

6. SCALING OF PHYSICS PROCESSES WITH LASER WAVELENGTH

The present section addresses our current understanding of how laser-plasma interactions vary with laser wavelength. This field is just beginning to be explored experimentally, particularly for laser wavelengths less than one micron. Therefore, it seems appropriate to outline the basic physics issues associated with wavelength scaling and to describe our preliminary theoretical predictions. We emphasize laser wavelengths of $1.06 \mu\text{m}$ and shorter, where frequency-multiplication of Nd laser light is expected to provide relevant experimental data.

There has already been a considerable amount of discussion [22,27] comparing CO_2 lasers ($\lambda_L = 10.6 \mu\text{m}$) with Nd-glass lasers ($\lambda_L = 1.06 \mu\text{m}$). By contrast, the behavior of plasma processes at wavelengths less than a micron raises new issues which have not yet been extensively investigated, and for this reason we shall emphasize submicron laser wavelengths.

The plasma physics coupling processes whose wavelength scaling we shall discuss are: 1) absorption of laser light, 2) hot electron generation and preheat, 3) stimulated scattering, 4) filamentation, and 5) transport inhibition. We shall show that virtually all of these plasma processes are expected to behave more favorably at very short laser wavelengths.

Of course the ultimate choice of wavelength for a fusion-reactor laser must also take into account the wavelength scaling of the basic hydrodynamic processes: ablation pressure, implosion efficiency, and implosion symmetry. We shall discuss these briefly, showing that ablation pressure and implosion efficiency improve dramatically for short-wavelength lasers. By contrast, good implosion symmetry is more difficult to achieve for short than for long laser wavelengths, and we discuss some of the design trade-offs which may allow improved symmetry for short-wavelength lasers.

In general, the results of this section emphasize the strong potential benefits of lasers with wavelengths considerably less than $1 \mu\text{m}$.

6.1 PLASMA PHYSICS AND LASER-PLASMA COUPLING

6.1.1 Absorption of Laser Light

It is thought that the most important absorption processes in present experiments are inverse Bremsstrahlung and resonance absorption. These were discussed in detail in Section 2. The two mechanisms dominate in different regions of parameter space, as shown schematically in Fig. 41. For high laser intensities I_L and long laser wavelengths λ_L , the underdense plasma tends to be collisionless and resonance absorption is strong. Absorption is roughly 20–30% independent of λ_L in this regime.

Inverse Bremsstrahlung (collisional absorption) dominates at short laser wavelengths and lower intensities. The total absorption fraction integrated over the underdense plasma is given by Eq. (20). For weak absorption, Eqs. (20) and (14) yield a total absorption which scales as $\nu_{ei}(n_e)L$, where

$$\nu_{ei}(n_e) \propto Z n_e T_e^{-3/2} \propto Z \lambda_L^{-2} T_e^{-3/2} \quad (100)$$

A simple argument [120] estimates the scale length as

$$L = c_s \tau_L \propto (Z T_e)^{1/2} \tau_L \quad (101)$$

and obtains the temperature T_e from the absorbed laser intensity:

$$\alpha_{\text{abs}} I_L = f n_e m_e (T_e/m_e)^{3/2} \quad (102)$$

FIG. 41. Regions of (laser intensity, laser wavelength) parameter space where collisional absorption and resonance absorption are dominant. Inverse Bremsstrahlung is strongest at short λ_L and low I_L ; resonance absorption at high values of $I_L \lambda_L^2$.

Equations (100)–(102) can be solved for T_e and α_{abs} , yielding

$$\alpha_{\text{abs}} \propto \frac{(Z^{3/2} \tau_L)^{0.6}}{I_L^{0.4} \lambda_L^2} \quad (103)$$

Thus inverse Bremsstrahlung is predicted to be stronger for high target Z , long laser pulses τ_L , low intensities I_L and especially for short laser wavelengths λ_L . These theoretical trends are confirmed in the data of Fig. 42, from experiments at École Polytechnique [121].

6.1.2 Hot Electron Generation and Preheat

Suprathermal electrons are produced when laser light is absorbed by resonance absorption at the critical density surface, and by the Raman and two-plasmon decay processes at densities $n_e < n_c/4$. These hot electrons can have long mean free paths, and can thus cause preheat of inner target layers or of the fuel region. Consequences of such preheat include decreased peak fuel densities, or else higher required laser energies to attain a given central density.

Recent research on resonance absorption has shown a somewhat weak dependence of the hot electron temperature T_h on laser wavelength; Eqs. (30) and (31) imply that

$$T_h \propto \lambda_L^{0.6-0.8} \quad (104)$$

The range of the hot electrons is proportional to T_h^2 , and thus

$$\text{hot electron range} \propto \lambda_L^{1.2-1.6}$$

The range is more sensitive to λ_L , and can be greatly reduced by the use of short-wavelength lasers.

Designs for reactor-sized targets are most vulnerable to preheat from electrons more energetic than the 10- to 50-keV electrons produced by resonance absorption. This is because the outer shell of the target tends to be on the order of 1/2 mm thick, and thus electrons with $E \ll 2$ MeV will not penetrate to the target

FIG. 42. Experimental measurements of absorption versus laser intensity from the École Polytechnique [121], for three different laser wavelengths. Short-pulse data are indicated by triangles at $\lambda_L = 1.06 \mu\text{m}$ and $0.53 \mu\text{m}$, and long-pulse data by circles. The two asterisks in the upper right-hand corner represent short-pulse absorption measurements at $\lambda_L = 0.26 \mu\text{m}$. The general trends follow those predicted by Eq. (103): absorption increases for shorter laser wavelengths, lower laser intensities, and longer laser pulse lengths.

interior. Hence the possible existence of small, very high-energy tails on the electron distribution becomes a crucial part of any preheat prediction.

Computer simulations with large underdense plasmas [48] ($n > 100 n_c$) have shown that the Raman and two-plasmon decay processes can sometimes produce very hot electron tails, which would be of potential concern for preheat in reactor targets. As a result, prudent target design suggests that one try to stay below the thresholds for the Raman and two-plasmon decay instabilities.

One can stay below the threshold due to inverse Bremsstrahlung absorption of the incoming light [48] (Section 3.2):

$$I_l \sim 10^{12} \frac{\text{W}}{\text{cm}^2} \times \begin{cases} \lambda_l \left(\frac{1 \mu\text{m}}{\lambda_l} \right)^3 \left(\frac{1 \text{ keV}}{T_e} \right)^{1/2} & \text{Inverse Bremsstrahlung} \\ \lambda_l^{3/4} \left(\frac{1 \mu\text{m}}{\lambda_l} \right)^{11/4} \left(\frac{1 \text{ keV}}{T_e} \right)^2 & \text{Raman absolute} \end{cases} \quad (108)$$

Alternatively, one can suppress the instability effects by density gradient alone [48]:

$$I_l \sim 10^{16} \frac{\text{W}}{\text{cm}^2} \left[\left(\frac{1 \text{ keV}}{T_e} \right) \left(\frac{1 \mu\text{m}}{\lambda_l} \right) \left(\frac{1 \mu\text{m}}{\lambda_l} \right) \right] & \text{Inverse Bremsstrahlung} \quad (109)$$

$$I_l \sim 10^{18} \frac{\text{W}}{\text{cm}^2} \left[\left(\frac{1 \mu\text{m}}{\lambda_l} \right)^{11/4} \left(\frac{1 \mu\text{m}}{\lambda_l} \right)^{2/4} \right] & \text{Raman absolute} \quad (110)$$

These expressions are derived for a linear density profile of length L from $n = 0$ to $n = n_c$. Both processes criteria are considerably easier to meet for shorter wavelength lasers, with λ_l a fraction of a μm .

A final preheat issue concerns corona-core decoupling [89,90]. Hot electrons tend to preferentially populate the outer regions of the corona, where they have very long mean free paths. Figure 34 illustrates this situation. A hot electron can bounce many times in the corona before it hits the core [89,90]. The number of bounces is proportional to $(R_{\text{corona}} - R_{\text{core}})/r_D$. If we estimate the outer radius of the corona by

$$R_{\text{corona}} \sim \tau_l / (I_h M_p)^{1/2} \quad (111)$$

and use the resonance absorption scaling $I_h \propto (I_l \lambda_l^2)^{1/2}$, we obtain [122] an energy loss to the corona proportional to $\tau_l (I_l \lambda_l^2)^{1/4}$. Thus one expects lower coronal losses for shorter laser wavelengths and shorter laser pulses.

6.1.3 Stimulated Scattering

The wavelength scaling of Raman scattering was discussed in the previous section, being directly related to preheat. Here we discuss stimulated Brillouin scattering, which can cause preheat, reflectivity, and decreased absorption of the laser light. Three different physical regimes will be discussed: 1) low density that the ion waves break, 2) intermediate values of $I_l \lambda_l^2$, where ion waves stay linear but the waves are collisionless, 3) the collisional limit where inverse Bremsstrahlung absorption is important. In each case, we consider a design for a given reactor-sized target, so that the plasma size remains fixed as λ_l varies.

1) Very High Values of $I_l \lambda_l^2$

Here the ion waves break, and they are nonlinearly limited to a maximum amplitude $\delta n_i \sim 10^{-3}$. The theory of Section 3 for this case gives an SBS gain

$$Q^{\text{NL}} \propto \left(\frac{L}{\lambda_l} \right) \left(\frac{\delta n_i}{n_i} \right)^{\text{NL}} \propto \lambda_l^{-1} L \quad (112)$$

This may be the relevant case for CO_2 lasers, and would act together with profile steepening to reduce SBS when $\lambda_l = 10.6 \mu\text{m}$. In this regime, SBS is reduced for longer-wavelength laser light.

2) Intermediate values of $I_L \lambda_L^2$.

When the ion waves remain linear, we showed in Section 3.1 that the total SBS gain can be parametrized by

$$Q \propto \frac{1}{\lambda_L} \left(\frac{n_e}{n_c} \right) \left(\frac{E_L^2}{8\pi n_e T_e} \right) \quad (109)$$

We estimate the density n_e in Q by the hot-electron density n_h :

$$n_e \approx n_h = \frac{I_{\text{abs}}}{m_p (T_h + m_p)^{3/2}} \quad (110)$$

for an absorbed intensity I_{abs} . Similarly, $E_L^2 / 8\pi n_e T_e \propto I_L \lambda_L^2 / T_h$, and we use the hot-electron temperature scaling from Eq. (30): $T_h \propto (I_L \lambda_L^2)^{1/2}$. Substitution of these results into Eq. (109) yields

$$Q \propto \lambda_L \quad (111)$$

Thus SBS is predicted to diminish at short laser wavelengths, in this intermediate regime of $I_L \lambda_L^2$.

3) Limitation of SBS by Inverse Bremsstrahlung

For laser and target parameters where inverse Bremsstrahlung heating becomes stronger than resonance absorption heating, the wavelength scaling of Brillouin scattering changes. The effective damping rate for the light is $\nu_L = \nu_e (n_i/n_e) T_e^2$, where $\nu_e(n_i) \propto n_i T_e^{-3/2} Z$ represents the electron-ion collision frequency at the critical density. For a linear density profile $n_e(x) = n_c(x - 1/n)$, the incident intensity is reduced by a factor e^{-1} when the light has propagated a distance

$$L_{\text{abs}} = \left[\frac{3eL_n^2}{\nu_e(n_c)} \right]^{1/3} \quad (112)$$

Here we assume $L_{\text{abs}} \ll L_n$ and neglect swelling for simplicity. The Brillouin reflectivity must occur in a length less than a few times L_{abs} if it is to cause reflection of a substantial fraction of the incident light.

To evaluate Eq. (111) we need to know the self-consistent temperature T_e . The scaling of temperature with wavelength can be determined by requiring that the electrons where the light has penetrated one absorption distance be able to carry off the absorbed laser intensity I_{abs} ,

$$I_{\text{abs}} = n_e \left(\frac{I_{\text{abs}}}{L_n} \right) m_p^{-1/2} T_e^{3/2} \quad (113)$$

Then since $L_{\text{abs}} \propto (I_L^2 n_e^{-3})^{1/3}$, this implies that $T_e \propto \lambda_L^{2/3}$.

The effective Q which describes the Brillouin reflectivity is, for $n_i \ll n_e$,

$$Q_{\text{th}} = \frac{\pi}{2} \left(\frac{\nu_0^2}{\nu_{\text{te}}^2} \right) \int_{\lambda=0}^{L_{\text{abs}}} \frac{\lambda}{L_n} \left(\frac{d\lambda}{\lambda_L} \right) = \frac{\pi}{4} \left(\frac{\nu_0^2}{\nu_{\text{te}}^2} \right) \left(\frac{L_{\text{abs}}^2}{L_n \lambda_L} \right) \quad (114)$$

Thus for a fixed value of L_n we find the strong wavelength scaling

$$Q_{\text{th}} \propto \lambda_L^{7/3} \quad (115)$$

in the regime where inverse Bremsstrahlung is the dominant heating process. In Section 3.1.4 we showed that this favorable scaling with short laser wavelengths holds for a wide variety of profiles [Eq. (54)]. For a uniform plasma [53] the above arguments yield the scaling $Q_{\text{th}} \propto \lambda_1^{-1/2}$.

In summary, for all but the highest-intensity regime where $E_1^2 8\pi n_e T_e \gg 1$, SBS is predicted to become considerably less severe for short-wavelength lasers. We recall that an alternative way to reduce SBS gain is by using a laser with large bandwidth, or a laser with many appropriately spaced lines [56,57].

6.1.4 Filamentation

Like SBS, filamentation has several different scaling regimes which depend on the laser intensity. At very high intensities ($E_1^2 8\pi n_e T_e \gg 1$) the density profile of the underdense plasma is very steep. Computer simulations indicate [43] that the underdense plasma has a density of about $n_e/10$. Thus, filamentation would need a very long length of plasma to grow significantly.

In the intermediate-intensity regime, the number of spatial gain lengths for filamentation is given by Q in Eq. (109). The predicted wavelength scaling of filamentation gain is similar to that of SBS with $T_i \approx ZT_e$. From Eq. (110) we then have $Q \propto \lambda_1$; there should be less filamentation at shorter laser wavelengths. If $T_i \ll ZT_e$, both particle-in-cell and fluid computer simulations [123] show that Brillouin sidescatter is strong enough to prevent filament formation, because the coherence of the filaments is destroyed by sidescatter. This suggests that if sidescatter were suppressed, for example, by broad laser bandwidth, filamentation might re-emerge.

In the collisional regime of low laser intensities and short wavelengths, filamentation in principle becomes dominated by the joule-heating nonlinearity rather than by the ponderomotive force (Perkins and Valeo [75], Sodha et al. [71]). The practical importance of thermal or collisional self-focusing for laser fusion is not yet clear [124].

6.1.5 Transport Inhibition

Inward thermal transport by electrons is expected to be less inhibited for short-wavelength lasers than for long-wavelength lasers.

Recent experiments suggest that thermal-electron heat flux is strongly inhibited below its classical value, $q_e = -\kappa T_e$. The mechanism producing inhibited transport has not yet been precisely identified, but it is generally attributed to a variety of plasma processes that produce "effective" collisions at a rate greater than the classical collision frequency, $\nu_{\text{eff}} > \nu_{ei}$. Since electron heat flow is a diffusive process, a greater collision rate ν_{eff} produces an inhibited velocity of heat flow that is slower than classical.

For the purposes of inward heat transport to dense parts of the target, we are interested in the values of ν_{ei} and ν_{eff} at an electron density several times the critical density. Since $\nu_{ei} \propto n_e T_e^{-3/2} \propto \lambda_1^{-2} T_e^{-3/2}$, we see that classical collisions increase rapidly with shorter laser wavelengths. As the laser wavelength becomes shorter, it then becomes more and more difficult for an "anomalous" plasma physics process to produce an effective collision frequency ν_{eff} comparable with ν_{ei} . Transport inhibition in the over-dense plasma should thus be less strong for short-wavelength lasers.

Consider magnetic fields as a case in point. Megagauss dc magnetic fields are thought to be generated by nonparallel density and temperature gradients. By balancing the magnetic-field source with loss due to fluid convection,

$$\frac{e}{en_e} \nabla n_e \times \nabla T_e \approx \nabla \times (\mathbf{V} \times \mathbf{B}) \quad , \quad (115)$$

we find by dimensional analysis that the product of the electron cyclotron frequency Ω_{ce} and the electron-ion collision time τ_{ei} is

$$\Omega_{ce} \tau_{ei} = \frac{\lambda_{mfp}}{L_T} \left(\frac{M_i}{Zm_e} \right)^{1/2} \quad , \quad (116)$$

where λ_{mfp} is the scattering mean free path of an electron; L_T is the temperature gradient scale length; and M_i , m_e are the ion and electron masses.

One expects significant inhibited transport when $\Omega_{ce}\tau_{ei} \gg 1$. The temperature gradient length L_T is at least as large as the impressed gradient in the laser-illumination pattern. On the other hand the mean free path λ_{mfp} at a few times the critical density scales as $n_c^{-1} T_e^2 \propto \lambda_L^2 T_e^2$. Thus short laser wavelengths tend to sharply decrease the value of $\Omega_{ce}\tau_{ei}$, and consequently decrease the inhibiting effect of magnetic fields on thermal transport into the over-dense regions.

In a second class of proposed mechanisms for transport inhibition [93,95], a crucial role is played by the collisionless outward streaming of suprathermal electrons produced at the critical surface. Electric fields which form to oppose the outward flux of suprathermals can have a variety of transport-inhibiting effects via their interaction with the thermal-electron population. However, short-wavelength lasers cause conditions where the suprathermal electrons are more collisional at the critical surface, since at the critical density the mean free path scales as

$$\lambda_{mfp}(n_c) \propto T_e^2/n_c \propto \lambda_L^{3.6} \quad (117)$$

This is a very strong scaling with wavelength. It suggests that lasers having wavelengths in the 1-2- to 1-4- μ m range might be able to eliminate entirely this class of transport inhibition mechanisms: "suprathermal" electrons born at the critical surface would be thermalized right in the critical region, rather than free-streaming into the underdense plasma and creating large electric fields there.

We conclude that many classes of mechanisms which inhibit the inward electron heat flux are expected to become weaker as the laser wavelength decreases below 1 μ m.

6.2 HYDRODYNAMICS AND IMPLSION PHYSICS

The emphasis of the present report has been mainly on laser-plasma interaction and coupling physics. However for the topic of wavelength scaling, it seems appropriate to include a brief outline of the effect of laser wavelength on implosion physics. The reader is cautioned that this is not intended to be a full review of the field.

We shall treat three areas of implosion physics: symmetry, ablation pressure, and hydrodynamic efficiency.

6.2.1 Implosion Symmetry

For a situation where many laser beams irradiate a spherical target, there will inevitably be some degree of nonuniformity in the distribution of absorbed laser energy. Such an impressed nonuniformity can lead to degraded target performance, particularly for shell targets with high convergence ratios.

Several years ago Theissen and Nuckolls [125] and Mead and Lindl [126] explored the circumstances under which laser illumination nonuniformities can be smoothed by heat conduction in a hot atmosphere surrounding the dense target. They found that initial temperature nonuniformities with spherical harmonic number l at the critical surface r_c were smoothed by the factor $(R_a/r_c)^l$, when they reached the ablation surface R_a . Mead and Lindl found that for practical cases of interest, an imposed nonuniformity of 10% could thus be tolerated if the critical radius r_c were larger than roughly twice R_a during the first part of the implosion.

A second requirement for conduction smoothing [125,126] is that the atmosphere must be hot enough for the thermal conduction time to be small compared with the sonic transit time. For a target with $Z = 10$, $A = 20$, this requirement becomes

$$T(r_c) \gg T_{min} = 4.8 \left(\frac{L}{1 \text{ cm}} \right)^{1/2} \left(\frac{1 \text{ } \mu\text{m}}{\lambda_L} \right) \text{keV} \quad (118)$$

Here L is the size of the atmosphere.

With the above two smoothing requirements in mind, we can use the simple spherical steady ablation model of Ref. [1] to ask about the wavelength-scaling of symmetrization by a hot atmosphere. (This model

does not apply to CO₂ lasers, where extreme profile steepening and hot electrons dominate coronal behavior, nor to strongly accelerated targets.) Reference [1] predicts the radius and temperature of the critical surface as

$$\frac{r_c}{R_a} = 1.6 \left(\frac{I_{abs}}{10^{14} \text{ W/cm}^2} \right)^{0.11} \left(\frac{\lambda_L}{1 \mu\text{m}} \right)^{0.38} \left(\frac{0.1 \text{ cm}}{ZR_a} \right)^{0.08} \quad (119)$$

$$\frac{T_c}{1 \text{ keV}} = 1.7 \left(\frac{I_{abs}}{10^{14} \text{ W/cm}^2} \right)^{0.52} \left(\frac{\lambda_L}{1 \mu\text{m}} \right)^{0.83} \left(\frac{ZR_a}{0.1 \text{ cm}} \right)^{0.11} \quad (120)$$

for a flux limit comparable to the experimentally inferred value $f = 0.03$. Here I_{abs} is the absorbed laser intensity evaluated at the surface of the dense target, R_a .

The first criterion for effective smoothing is that $r_c/R_a \geq 1.5$ –2. Consider for example a target with ablation radius $R_a = 0.1 \text{ cm}$, $Z = 10$, $A = 20$, and $I_{abs} = 10^{14} \text{ W/cm}^2$. Then Fig. 43a shows the ratio of the critical to ablation radii r_c/R_a as a function of laser wavelength λ_L (Eq. 119). It is clear that in this example, one needs a laser wavelength $\lambda_L \geq 2$ –3 μm to be in the desired smoothing regime $r_c \geq 1.5$ –2 R_a .

The second criterion for good smoothing is that the temperature at critical, given by Eq. (120), be large compared with the minimum temperature T_{min} derived from conduction-time arguments in Eq. (118). Figure 43b shows that for $\lambda_L > 1 \mu\text{m}$, the coronal temperature T_c predicted by the model in Ref. [1] is indeed large compared with T_{min} .

Thus the results using the simple ablation model of Ref. [1] are in substantial agreement with the more detailed numerical analysis of Mead and Lindl: each of the two smoothing criteria are best met with lasers of relatively long wavelength, $\lambda_L > 2 \mu\text{m}$, if one wants to use initial nonuniformities of $\sim 10\%$ in the absorbed laser energy.

In the first part of this wavelength-scaling section, we showed that almost all of the laser-plasma coupling issues were improved for laser wavelengths considerably smaller than $1 \mu\text{m}$. In the discussion to follow we shall show that ablation pressure and hydrodynamic efficiency are also higher for short-wavelength lasers. Thus the symmetry issue seems to be the only one where long wavelengths, $\lambda_L > 2 \mu\text{m}$, are more desirable. One can think of several scenarios for compromise between the conflicting requirements of long laser wavelength for symmetry and short wavelength for the remainder of the target physics:

a) References [125] and [126] suggested using a prepulse with $\lambda_L \approx 2 \mu\text{m}$ to create the hot, smoothing atmosphere, followed by a main shorter-wavelength laser pulse to do the actual implosion. This might allow the best of both worlds, since uniformity requirements could be relaxed, yet one could still use the optimum short-wavelength laser for target compression.

b) One could use a short-wavelength laser for the entire heating and compression process.

Under these conditions there would not be a hot, smoothing atmosphere. Consequently, the required symmetry of laser light absorption would be considerably higher. Such a scenario may be feasible by using a larger number of laser beams, and by taking advantage of an appropriate mix of inverse Bremsstrahlung and resonance absorption to optimize the focusing scheme [127].

Both Fabre [128] and Theissen and Nuckolls [125] have noted that these exacting symmetry requirements might be relaxed if one used smaller-radius targets at the shorter laser wavelengths. For example, one could keep plasma coupling processes well in control by using a constant value of $I_L \lambda_L^2$, and increasing the laser intensity I_L as one decreased λ_L . The resulting smaller targets could have lower aspect ratios and perhaps lower convergence ratios. As a result, one might be able to tolerate nonuniformities larger than the 1–2% level derived by Mead and Lindl for a high convergence-ratio target. A smaller-radius target might also allow greater flexibility in choice of a focusing scheme for optimum uniformity of absorption, because the ratio of laser beam size to target size could be larger. This option of using smaller-radius targets for shorter laser wavelengths needs further quantitative evaluation from the target-design point of view.

c) A third, quite different way to solve the symmetry problem is to convert the incident short-wavelength laser light to x rays and then to implode the target using x rays. The latter approach has the potential of significantly relaxing requirements of illumination uniformity. One is then able to reap the advantages of short-wavelength light for coupling physics, without impairing implosion symmetry. Using this approach,

FIG. 43. Wavelength-scaling predictions of a simple spherical ablation model [1]. For this example the ablation radius is $R = 0.1 \text{ cm}$, the absorbed laser intensity (normalized to R_a) is $I_{\text{abs}} = 10^{14} \text{ W/cm}^2$, and the target material has $A = 20$, $Z = 10$. a) Variation of critical-surface radius r_c with laser wavelength λ_L . b) Variation of critical-surface electron temperature T_c and minimum smoothing temperature $T_{\min}$, Eq. (118), with laser wavelength λ_L . For $\lambda_L > 2 \mu\text{m}$, Mead and Lindl's two symmetry conditions [126] are roughly fulfilled.

recent experiments at the Lawrence Livermore National Laboratory have imploded targets to fuel densities between 100 and 200 times the density of liquid D-T (0.2 g cm^{-3}), using $1.06\text{-}\mu\text{m}$ laser light.

To summarize the results of this symmetry discussion, optimum smoothing by a large, hot atmosphere is obtained when a relatively long laser wavelength is used to create that atmosphere ($\lambda_L > 2 \mu\text{m}$). Since other target physics processes optimize at wavelengths considerably shorter than $2 \mu\text{m}$, one might want to use a long-wavelength prepulse to create the atmosphere, followed by a short-wavelength main pulse. Alternatively one might be able to achieve adequate symmetry with a short-wavelength laser alone, by using more laser beams and smaller targets, or by converting the laser energy to x-ray energy before doing the implosion. Since the symmetry issue is of the essence for high-gain, high convergence-ratio targets, it is crucial to develop a more quantitative understanding of the advantages and disadvantages of these different alternatives.

6.2.2 Ablation Pressure

Inward motion of the dense imploding target material is driven by the ablation pressure of the heated outflowing plasma. One must therefore understand how ablation pressure scales with laser parameters such as intensity and wavelength. The main effects may be illustrated heuristically using the simplified spherical ablation model of Ref. [1]. (Recall that this model is not valid for a CO_2 laser-driven implosion, or for strong inward acceleration.)

According to Ref. [1], the ablation pressure p_a and ablation rate $\dot{m}$ are given by

$$p_a = 7.2 \left[\left(\frac{I_{\text{abs}}}{10^{14} \text{ W/cm}^2} \right)^{0.57} \left(\frac{1 \mu\text{m}}{\lambda_L} \right)^{0.99} \left(\frac{Z R_a}{0.1 \text{ cm}} \right)^{0.07} \right] \text{ Mbar} \quad (121)$$

$$\dot{m} = 2.3 \times 10^{-4} \left[\left(\frac{I_{\text{abs}}}{10^{14} \text{ W/cm}^2} \right)^{0.48} \left(\frac{1 \mu\text{m}}{\lambda_L} \right)^{0.82} \left(\frac{R_a}{0.1 \text{ cm}} \right)^{1.89} \right] \text{ g/s} \quad (122)$$

for a flux limit $f \approx 0.03$. Thus for example, the ablation pressure scales roughly as

$$p_a \propto I_{\text{abs}}^{0.6} / \lambda_L \quad (123)$$

Shorter wavelengths produce higher ablation pressures. Figure 44 illustrates the change in pressure profile with laser wavelength, for $f \approx 0.06$.

FIG. 44. Pressure profiles predicted by simple ablation model for three different laser wavelengths, with the same absorbed laser power $P_L = 5 \times 10^{14} \text{ W}$ and target parameters $R_a = 0.07 \text{ cm}$, $Z = 6$, $A = 12$, and flux limit $f = 0.06$. Shorter laser wavelengths produce higher peak pressures at $r = R_a$. Similar behavior was predicted in Refs. [129] and [130] for classical heat flow.

To assess the impact of this scaling, consider a *given* target design, optimized to perform at a fixed value of the ablation pressure p_a . Equation (123) means that if a longer-wavelength laser is used to drive this target, a higher laser intensity will be required to obtain the required fixed ablation pressure p_a . Equation (123) says that necessary driving intensity scales as

$$I_L \propto (p_a \lambda_L)^{5/3} \quad (124)$$

Thus an increase of a factor of 2 in the laser wavelength will require 3.2 times the laser intensity to drive the same target design. As a result $I_L \lambda_L^2$, the parameter which characterizes many laser-plasma coupling processes, will increase by a total factor of 13 for each doubling of λ_L . Arguments such as this make short-wavelength lasers look very promising indeed.

Similar trends hold for the mass ablation rate, $\dot{m}$. For spherical targets, Eq. (122) predicts roughly

$$\dot{m} \propto I_{abs}^{1/2} / \lambda_L^{0.8} \quad (125)$$

Although this prediction has not yet been directly tested in spherical geometry, a recent plane-geometry experiment has confirmed that the mass ablation rate increases dramatically at short laser wavelengths. Fabbro et al. [79] measured the amount of ablated material passing outward through the critical surface in a given amount of time. This should be proportional to the ablation rate $\dot{m}$. The experiment used three different techniques to determine the depth of the heated region as a function of wavelength. Figure 45 shows one of these measures: the foil thickness producing 10% transmission of the incident laser light. (Foil material was CH, and the data are normalized to $I_L = 10^{15}$ W/cm², $\tau_L = 100$ ps for comparison purposes.) It is obvious from Fig. 45 that much more matter is ablated away for $\lambda_L = 1/4 \mu\text{m}$ than for $\lambda_L = 1 \mu\text{m}$.

6.2.3 Hydrodynamic Efficiency

By implosion or hydrodynamic efficiency we mean the inward kinetic energy of the imploding target, relative to the total absorbed laser energy. With a given amount of inward electron heat flux, newly heated material at the surface of the cold dense target expands outward. By conservation of momentum, the cold remainder of the target then gets a push inward, due to the rocket effect. Figure 29 illustrates this process.

FIG. 45. Experimental measurement by Fabbro et al. [79] of the foil target thickness producing 10% transmission of laser light, at three different laser wavelengths. Data are normalized to $I_L = 10^{15}$ W/cm², $\tau_L = 100$ ps for comparison purposes. Foil material was CH. It is clear that much more matter was heated for $\lambda_L = 1/4 \mu\text{m}$ than for $\lambda_L = 1 \mu\text{m}$.

Transport inhibition influences the implosion mechanism because as inward transport becomes impeded, smaller and smaller fractions of the absorbed laser intensity are put into inward heat flow. Instead, the absorbed intensity goes increasingly into driving the expansion of plasma outside the critical surface. Thus, the inner parts of the target receive less energy flux and implode less vigorously.

For cases where the ponderomotive force is weak, the critical density surface is farther from the ablation surface for long laser wavelengths than for short ones. This is illustrated in Eq. (119). As a consequence, we shall show that transport inhibition has a relatively larger impact on implosion properties for longer-wavelength lasers, since the inward-diffusing electrons have to transport their energy over a longer distance.

The simple spherical ablation model of Ref. [1] can be extended to derive the hydrodynamic efficiency of an implosion [55]. One starts by integrating the momentum equation

$$\rho \frac{dv_r}{dt} = - \frac{\partial p}{\partial r} \quad (126)$$

over a spherical volume whose outer radius is just inside the ablation radius R_a . When the shell moves inward together, one obtains

$$m(t) \frac{d\bar{v}}{dt} = - 4\pi R_a^2 p_a \quad (127)$$

where $m(t)$ is the mass remaining in the shell at time t , and $\bar{v}$ is the mean inward velocity of the shell.

The ablation theory of Ref. [1] gives expressions for the scaling of the mass ablation rate $\dot{m}$ and the ablation pressure p_a with the shell radius R_a :

$$\dot{m} = \dot{m}_0 (R_a/R_0)^a, \quad p_a = p_0 (R_a/R_0)^{-b} \quad (128)$$

Here R_0 , $\dot{m}_0$, and p_0 are the initial values of R_a , $\dot{m}$, and p_a . The exponents a and b in Eq. (128) depend only on the flux limit f . One can now divide Eq. (127) by $\dot{m}$, to obtain

$$m \frac{d\bar{v}}{dm} = \left(\frac{4\pi R_0^2 p_0}{\dot{m}_0} \right) \left(\frac{R_0}{R_a} \right)^{b+a-2} \quad (129)$$

Numerically, the value of the exponent $b + a - 2$ is quite close to zero. For the experimentally observed value of the flux limit $f = 0.03$, one has $b + a - 2 = 10^{-2}$, and $b + a - 2$ remains less than 10^{-1} for all flux limits $f \leq 0.6$. We shall make the approximation $b + a - 2 = 0$ in Eq. (129), and integrate to obtain the rocket equation

$$\bar{v} = \left(\frac{4\pi R_0^2 p_0}{\dot{m}_0} \right) t \ln \left[\frac{m(t)}{m_0} \right] \quad (130)$$

Thus in spherical geometry, including the effects of the full temperature, density, and velocity profiles, one obtains a rocket equation exactly analogous to the usual planar "cold rocket" result

$$\bar{v}(t) = v_{ex} t \ln \left[\frac{m(t)}{m_0} \right] \quad (131)$$

In the spherical case the effective exhaust velocity is

$$v_{ex} = \frac{4\pi R_0^2 p_0}{\dot{m}_0} \quad (132)$$

with the initial ablation pressure p_0 and mass ablation rate $\dot{m}_0$ evaluated using expressions (121) and (122) (or their analogues for different values of the flux limit $f = 0.03$).

Equation (130) for the mean inward velocity of the shell can now be used to find the hydrodynamic efficiency ϵ :

$$\epsilon = \frac{(1/2) m(t) \bar{v}^2}{P_{\text{abs}} \tau_L} = \left[\frac{2\pi R_0^2 p_0^2}{I_{\text{abs}} \dot{m}_0} \right] \left(\frac{x \ell n^2 x}{1-x} \right), \quad (133)$$

where P_{abs} and I_{abs} are the absorbed laser power and intensity, τ_L is the laser pulse length, and $x = m(t)/m_0$ is the fraction of mass remaining in the shell after a time t .

The factor in parentheses in Eq. (133) optimizes for $x = 0.2$, when 80% of the initial shell mass has been ablated away. We can obtain the wavelength scaling of the hydrodynamic efficiency ϵ by examining the factor in square brackets in (133). Using the results of Ref. [1] for p_0 and $\dot{m}_0$ as a function of flux limit f and laser wavelength λ_L , it is straightforward to show that

$$\epsilon \propto \lambda_L^{-p}, \quad (134)$$

where the exponent p varies from 1.2 to 0.5, as the flux limit f changes from 0.03 to 0.4. Figure 46 shows the exponent p as a function of f .

Within the validity limits of the simple ablation model (few hot electrons, weak ponderomotive force, gentle acceleration), it is clear from Eq. (134) and Fig. 46 that implosion efficiency increases dramatically for short-wavelength lasers. The increase is strongest for very inhibited transport, as explained at the beginning of this section. But even for nearly classical values of the flux limit, the implosion efficiency scales as $1/\sqrt{\lambda_L}$. For the case when transport is completely classical, Refs. [129] and [130] found similar results: the implosion efficiency is highest at the shortest laser wavelengths.

FIG. 46. The simple spherical ablation model of Ref. [1] predicts that hydrodynamic implosion efficiency ϵ is proportional [55] to λ_L^{-p} , where λ_L is the laser wavelength. This figure plots the exponent "p" as a function of the flux limit f , which measures the degree of transport inhibition. In all cases the efficiency is best for short laser wavelengths.

The results of our wavelength-scaling discussion may be summarized as follows. If we consider laser wavelengths between a few microns and a fraction of a micron, we have shown that laser-plasma coupling, ablation pressure, and implosion efficiency are all more favorable for the shortest laser wavelengths. By contrast, implosion symmetry seems easier to achieve for longer-wavelength lasers. We discussed several ways in which symmetry might be improved for short-wavelength lasers. If these methods for better symmetrization prove effective, the results of this section point strongly to the potential benefits of submicron laser wavelengths.

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