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Indications of Strongly Flux-Limited Electron Thermal Conduction in Laser-Target Experiments*

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(Received 9 December 1974)

It is shown by comparison with calculations that anomalies in the results of intense laser irradiation of solid targets, including two-humped ion distributions, indicate a reduction of electron thermal conduction to considerably below classical values. This reduction is interpreted as a flux limit and appears to be sufficiently restrictive to modify significantly the design of laser fusion targets.

Most of the recent interest in laser-produced plasmas has focused on conditions under which electron thermal conduction is expected to play a central role in the transport of energy in the plasma. In particular, in laser fusion applications, the ablation-driven compression and thermal heating of the target core are strongly influenced by the manner in which the laser energy deposited near the critical density ($n_c = 10^{21}$ cm^{-3} for Nd-glass and $n_c = 10^{19}$ cm^{-3} for CO_2 lasers) is transported into the target interior. It is the purpose of this Letter to point out that certain anomalies seen recently in high-power-density laser-plasma experiments indicate that electron thermal fluxes near critical density are limited to much less than the classical or mildly flux-limited values that have previously been assumed. This more stringent flux limiting would constitute an energy decoupling mechanism which would significantly modify design of fusion targets.

Previous predictions of experimental results were based on calculations of coupled hydrodynamics and heat flow in which the electron thermal flux was given by

$$\begin{aligned} F &= \begin{cases} F_c, & F_c < F_l, \\ F_l, & F_l < F_c, \end{cases} \\ F_c &= -K_c \nabla T_e, \\ F_l &= f n_e k T_e (k T_e / m_e)^{1/2}, \end{aligned} \tag{1}$$

where K_c is the classical conductivity of Landshoff and Spitzer[1] and F_l is an upper limit imposed to approximate the free-streaming, high-flux failure of the perturbation derivation of K_c.[2] Here m_e, n_e, and T_e are the electron mass, number density, and temperature, respectively, and f is a dimensionless number which is calculated to be between 0.5 and 1.0, depending on assumptions about collisionless processes, but is commonly taken to be about 0.6.[3,4] Here we show, by detailed experimental-calculational comparisons, that $f = 0.6$ is highly inconsistent with experimental results, but that the discrepancies are largely removed and certain new experimentally observed effects are given a plausible physical explanation by using the same form, Eq. (1), and taking $0.03 \lesssim f \lesssim 0.1$.

Since these values of f are so much smaller than those expected from free streaming, the observation that they are dictated by experiment is equivalent to identifying the presence of some microscopic process which is qualitatively new in this context. Investigation of the various possible processes is a major continuing effort and beyond the scope of this Letter. Here we only indicate several possibilities. Forslund[5] has shown that when F_c exceeds F_l with $f \simeq 0.1$ and $T_i/T_e \ll 1$, the plasma is unstable to the growth of ion-acoustic waves. The fluctuating electrostatic fields produced by the ion waves will inhibit the transport of energy by electrons. The form of Eq. (1) with $f \ll 1$ would, therefore, be consistent with ion-acoustic turbulence preventing heat flow from significantly exceeding the instability threshold. In principle, f would then be a function of T_i/T_e although if $T_i/T_e \ll 1$ the dependence should not be strong. Preliminary analysis and some two-dimensional numerical simulations indicate that the reduction of thermal conduction may be sufficient to account for the experimental observations.[6] However, adequate simulations are very costly and proceed slowly. Alternatively, internally generated magnetic fields[7] or electron turbulence[8] may reduce electron heat conduction.

The recent Los Alamos Scientific Laboratory experimental results which indicate reduced flux and with which we compare our phenomenological model are the following[9]: (1) The fraction of incident laser light transmitted by thin plastic foils is far less than predicted with $f = 0.6$. (2) The current of ions expanding from thick (30–150 μm) plastic foils shows a two-humped distribution as a function of velocity. The experiments indicate that an anomalously large fraction of the observed laser energy is converted to kinetic energy of the fast part of the ion distribution. Equation (1) with no flux limit ($f = \infty$) or a weak limit ($f \simeq 0.6$) does not predict either a two-humped distribution or the high ion energies observed. (3) The observed bremsstrahlung spectrum[10] for thick plastic foils has a higher intensity at photon energies near 50 keV than can be explained with classical electron conduction, if a Maxwell-Boltzmann electron distribution is assumed.

In this paper we consider only the simplest case, f = const, and we have replaced the expression for F of Eq. (1) with a harmonic mean $F = (F_c^{-1} + F_l^{-1})^{-1}$. The calculations were made with a one-dimensional Lagrangian hydrodynamics code with separate electron and ion temperatures and a perfect-gas equation of state. Laser energy is absorbed in the underdense ($n_e \lesssim n_c$) plasma very near the critical surface; absorption is neglected when the entire target becomes underdense. Except where noted otherwise, the laser pulse has a Gaussian time dependence cut off at the 1% level, a full width at half-maximum of 25 psec, a maximum intensity $\sim 5 \times 10^{15}$ W/cm^2, and a wavelength of 1.06 μm. These correspond approximately to the shape and duration of the pulse of the neodymium glass laser used in the experiments discussed herein. The laser energy was in the range 4 to 40 J and had a spot diameter at the target of about 100 μm. Such a pulse produces a plasma which is coin shaped and nearly one-dimensional during the time interval of laser energy deposition.

Figure 1(a) shows the calculated bremsstrahlung spectra for values of f between 0.01 and ∞ (no flux limit) superimposed on the experimental points of Ref. 10. Both the flux-limit parameter, f, and the laser energy were adjusted to obtain a match: Basically f controls the curvature of I_ν while the laser energy controls the magnitude of I_ν. Values of f between 0.03 and 0.1 give the correct curvature, while the fit with $f = 0.6$ to ∞ is very poor. With the experimental spot diameter of 100 μm[11] (following the publication of Ref. 10, it was determined that the laser spot diameter is closer to 100 μm than 50, as given therein[11]), the laser energy had to be reduced from the 10 J of Ref. 10 to 2 J to give the observed magnitude for I_ν. A reasonable fit could have been obtained with 3 or 4 J except at $h\nu < 20$ keV where the calculated values would be too large by a factor of about 3 to 5.

Stringent flux limiting ($f \lesssim 0.1$) causes the development of two thermal fronts, instead of the usual single front, as shown in Fig. 1(b). The front at n_c is stationary with respect to the critical surface and appears because the limited electron conduction flux is incapable of transferring all of the energy to the second front, which propagates into the target, as rapidly as the laser deposits energy in the subcritical blowoff. Note that in Fig. 1(b), the ions in the blowoff are very cold as the result of weak electron-ion coupling and expansional cooling. This does not prevent them from attaining large directed velocities as a result of the large electron pressure gradient.

The fraction of laser light incident on thin plastic foils which is transmitted when the foil goes underdense during the pulse is shown in Fig. 1(c) as a function of foil thickness. Again, $f \geq 0.6$

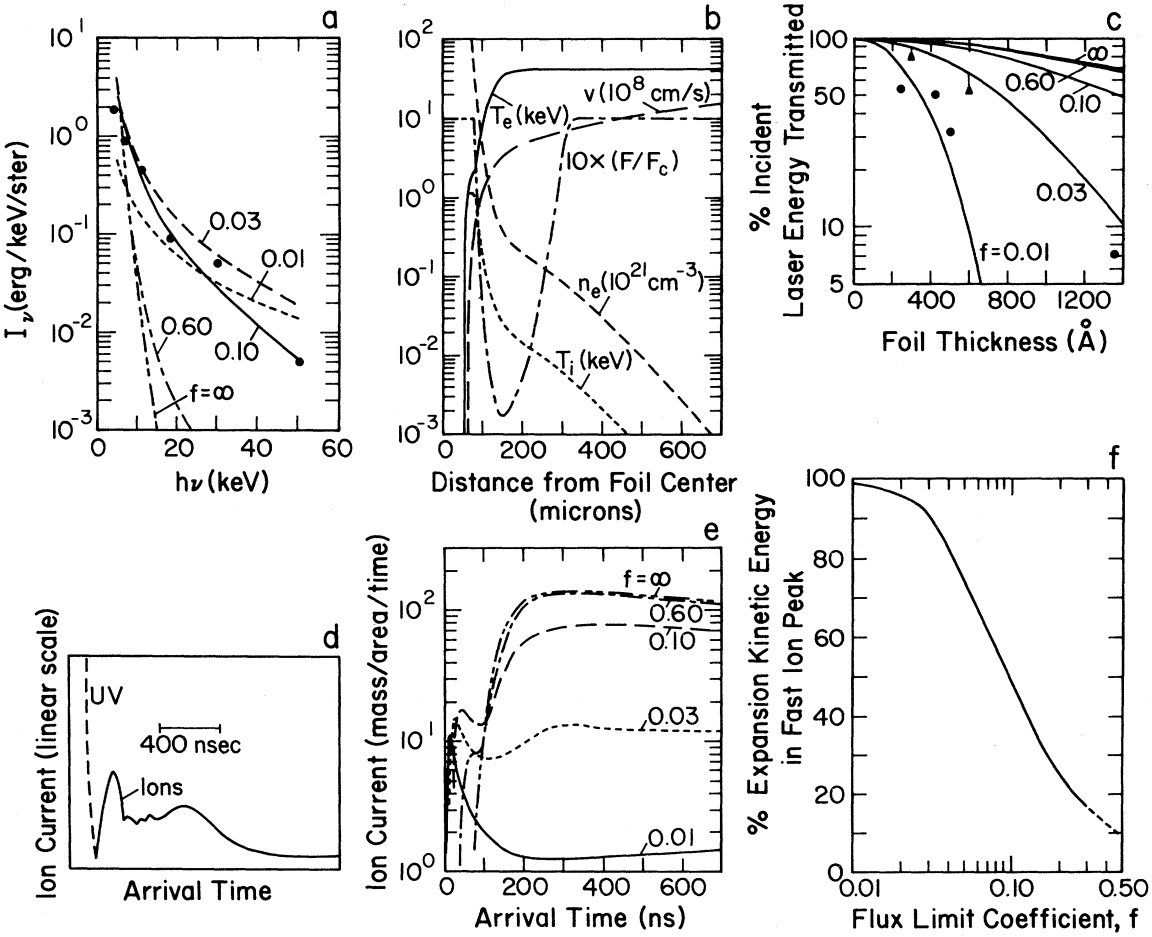

FIG. 1. (a) Calculated bremsstrahlung spectra I_ν versus photon energy $h\nu$ for different values of the flux limit coefficient f, compared to the experimental data (•) of Ref. 10. (b) Electron (T_e) and ion (T_i) temperatures, electron density (n_e), material velocity (v), and the ratio of F to F_c [see Eq. (1)] in a 150-μm CH_2 target irradiated with a 25-psec full width at half-maximum Gaussian pulse of maximum intensity 2×10^{16} W/cm^2; $f=0.1$ and the time shown is 60 psec. (c) The percentage of the incident laser energy transmitted by thin CH foils for different values of f using the experimental laser parameters (•, experimental data, Ref. 11; for the meaning of the ▲ points, see text). (d) Experimental ion current from a thick CH_2 foil irradiated by a Nd-glass laser (Ref. 11). (e) Calculated ion current with typical laser parameters (see text). (f) Calculated percentage of expansion ion kinetic energy residing in the fast ion peak versus flux limit coefficient f.

compares very poorly with experiment, while $f\simeq 0.01$ to 0.03 is much better. Because this is a dynamical effect, the results are much more sensitive to details of the spatial and temporal dependence of the experimental laser pulse and the true laser absorption mechanism. The former are known only crudely while the latter is perhaps not yet known at all. The two points in Fig. 1(c) denoted by triangles indicate the change from the $f=0.03$ curve when a parabolic laser pulse of 25 psec full width at half-maximum is used; because it has a shorter time of rise and fall than a Gaussian of the same width, more energy is forced into the target before hydrodynamic expansion makes the entire target underdense.

Finally, we consider the current of ions which are blown off of thick CH_2 foils. Figure 1(d) is an experimental oscilloscope trace showing an initial burst of uv radiation, followed by a sequence of fast-ion peaks, and finally the broad hump of slow plasma containing most of the heated target material. The fast-ion peaks appear to correspond to different ion charge-to-mass ratios, the separation being due to collisionless electrostatic acceleration in the blowoff region. Because the code used can only describe a single fluid, the charge separation effects could not be simulated. However, by reducing the flux limit coefficient, f, it was possible to produce a two-humped ion distribution where none existed when $f \gtrsim 0.6$ [see Fig. 1(e)]. Figures 1(d) and 1(e) are not directly comparable since Fig. 1(e) represents mass current, while Fig. 1(d) shows charge current, including secondary electrons knocked off the ion collector. But as Fig. 1(e) indicates, a distinct fast-ion peak only appears for $f \lesssim 0.1$.

For a fixed value of f, the kinetic energy ϵ of a proton at the point of maximum current of the fast-ion peak varies almost precisely as the $\frac{2}{3}$ power of I_{max}. This dependence follows from the removal of energy by ions from the absorption region near the critical density at a rate proportional to the incident laser flux: $n_c \epsilon^{3/2} \propto I$. In Fig. 1(f) we show the fraction of expansion kinetic energy contained within the fast-ion peak. It varies from nearly 100% at $f = 0.01$ to 15% at $f = 0.3$ where a break between the fast ions and warm plasma was barely discernible. Preliminary estimates from experiments[11] suggest that in the direction normal to the target, about 50% of the energy in ions appears as fast-ion kinetic energy, which is consistent with $f \simeq 0.1$.

We conclude that there is strong experimental evidence that some process is occurring in the target plasma which severely limits the electron conduction flux. Equation (1) with $f \simeq 0.03$ to 0.1 gives a good phenomenological agreement with at least three disparate experimental diagnostics of the plasma behavior. It is known that such a restrictive flux limit degrades the performance of laser-imploded spherical targets.[3] Also, recent experiments[12] and the form of F_l indicate that the thermal decoupling due to the flux limit is more severe at longer wavelengths where more powerful and efficient lasers are available. For these reasons it is of great importance to laser fusion research to obtain a better understanding of this aspect of the plasma behavior and of possible ways to modify or circumvent it.

*Work performed under the auspices of the U. S. Atomic Energy Commission.

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Stability of Envelope Solitons*

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One-dimensional envelope solitons are unstable against modulation in the direction perpendicular to their propagation for all amplitudes and wave numbers. This is a special case of the filamentation instability.

Large-amplitude plasma waves are known to be unstable against modulational instabilities.[1] The final state of the wave is assumed to be a set of envelope solitons, i.e. self-contained, stable, propagating wave packets of electron oscillations associated with a well in plasma density due to the ponderomotive force created by the oscillations. The k spectrum observed in one-dimensional computer simulations has been interpreted as arising from these solitons.[2] These