Efficient Light Absorption by Ion-Acoustic Fluctuations in Laser-Produced Plasmas*

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We describe a new efficient absorption mechanism for intense laser light in plasma resulting from an ion acoustic instability generated by a return current. Comparisons with experimental results are presented.

This Letter presents a new absorption mechanism which gives efficient absorption of intense light by a laser-produced plasma. The basic idea is that in absorption, the laser energy flux is converted into an electron thermal energy flux Q flowing into the plasma. In order for charge neutrality to be maintained, there must be a return current of low-velocity electrons flowing toward the laser (in the negative x direction). This return current excites ion acoustic waves, also propagating toward the laser. The laser light then experiences enhanced collisional damping on these ion density fluctuations in the underdense plasma. The absorption of the laser light then creates that very electron thermal energy flux which was required in the first place.[1-3] Values of Q/nmv_e^3 are small enough that a fluid model (dominated by anomalous transport) is valid.

Results of our theory seem to be in good qualitative agreement with many absorption, scattering, and x-ray measurements at Naval Research Laboratory (NRL).[4-7] It would be very difficult to explain these measurements by resonant absorption.[8] The measurements consistently show high fractional absorption (in excess of 50%) which is relatively independent of both polarization and angle of incidence.[5] In addition, NRL experiments indicate a fairly smooth critical surface for distance scales above about 1 μm.[7] Also, light absorption by enhanced ion density fluctuations would not tend to strongly produce nonthermal electrons. Energy flux is found to be carried principally by electrons at about two or three times the thermal speed. This also seems to be in agreement with hard-x-ray measurements.[4,6] Finally, layered-target experiments[4] indicate a value of Q/nmv_e^3 less than 0.2 which is also in agreement with our calculations. We now describe our calculations, and will close by giving more detailed comparisons with experimental results. The relevant steady-state fluid equations have been written out and discussed elsewhere.[1] We summarize them here as follows:

$$\frac{\partial}{\partial x}nT_e + neE + \frac{1}{4}\frac{e^2}{m^2\Omega^2}\frac{\partial}{\partial x}(E_i^2+E_r^2) = C_{ve}, \tag{1a}$$

$$\frac{\partial Q}{\partial x} = C_{Te} + \nu_{\rm an}\left(\frac{E_i^2+E_r^2}{8\pi}\right)\frac{\omega_{pe}^2}{\Omega^2}, \tag{1b}$$

$$\frac{3}{2}\frac{\partial}{\partial x}nT_e^2 + \frac{3}{2}\frac{T_e}{m}\left(neE + \frac{1}{4}\frac{e^2}{m^2\Omega^2}\frac{\partial}{\partial x}[E_i^2+E_r^2]\right) = C_{Qe}, \tag{1c}$$

$$nv = \text{const}, \tag{1d}$$

$$nMv\frac{\partial v}{\partial x} + \frac{\partial}{\partial x}nT_i - neE = -C_{ve}, \tag{1e}$$

$$\frac{\partial}{\partial x}\frac{3}{2}nvT_i + nT_i\frac{\partial v}{\partial x} = -C_{Te}, \tag{1f}$$

where E is the ambipolar electric field. $E_{i(r)}^2$ is square of the electric field of the incident (reflected) laser light, Q is the electron thermal energy flux in the x direction, Ω is the laser light frequency, $\nu_{\rm an}$ is the anomalous collision frequency, and C_{ve}, C_{Te}, C_{Qe} are quasilinear collision terms which describe the electron momentum, thermal energy and, thermal energy flux loss due to interactions with unstable waves.[1] In Eqs. (1a) and (1c), the effects of ponderomotive force have been included. All other notation is standard. A fluid description is valid as long as the gradient scale length is long compared to the anomalously reduced mean free path,[1] $V_e(e\varphi/T_e)^{-2}/\omega_{pe}$—a condition easily satisfied. Here φ is the fluctuating potential due to the instability and V_e is the electron thermal velocity.

Coupled to this fluid system are equations for $E_{i(r)}^2$ and also equations for the unstable ion acoustic waves. These are

$$c\frac{d}{dx}\left(\cos^2\theta - \frac{\omega_{pe}^2}{\Omega^2}\right)^{1/2}E_{i(r)}^2 = -(+)\nu_{\rm an}\frac{\omega_{pe}^2}{\Omega^2}E_{i(r)}^2, \tag{2}$$

where θ is the angle between the wave vector of the laser light and the x axis and

$$\frac{d}{dx}\left|\frac{e\varphi(k)}{T_e}\right|^2 = 2\left(\frac{\gamma}{v-(T_e/m)^{1/2}}\right)\left|\frac{e\varphi(k)}{T_e}\right|^2, \qquad (3)$$

where γ is the growth rate,[1] and k is the wave number of the unstable wave.

The final quantities to specify are C_{ve}, C_{Te}, C_{Qe}, and ν_{an}. The quantity ν_{an} depends on the component of $\vec{k}$ in the direction of $\vec{E}_i$ (for instance the y direction). Thus the angular spectrum of the ion acoustic fluctuations is needed. We make use of results of many numerical simulations of ion acoustic turbulence in two dimensions,[9] which show a cone of unstable waves out to an angle of between about 45 and 60 deg. We use this basic result and assume a three-dimensional conical spectrum uniform in angle up to 55° to the x axis and then dropping sharply to zero. Making this assumption, we find[1]

$$\nu_{an} = \Omega\frac{\omega_{pe}^2}{\Omega^2}\frac{3k_{min}}{k_{max}}\sum_{k=k_{min}}^{k_{max}}\left(\frac{k}{k_{max}}\right)^2 0.3\pi\left(\frac{\omega_{pe}}{\Omega}\right)^3 k\lambda_D\left[\left(1-\frac{\omega_{pe}^2}{\Omega^2}\right)^2+4\left(\frac{\omega_{pe}^2}{\Omega}\right)^6 k^2\lambda_D^2\right]^{-1}\left|\frac{e\varphi(k)}{T_e}\right|^2, \qquad (4)$$

where k_{min} and k_{max} are the minimum and maximum wave numbers included in the summation. In our calculations we take $k_{max}=0.75k_D$ and $k_{min}=k_{max}/9$. The C's are calculated as in Ref. 1, only making the same conical approximation to the wave spectrum as made in the calculation of ν_{an}.

Now it is worthwhile pointing out that if a transverse magnetic field exists, as is usually the case for laser light focused on a slab,[10] the ion acoustic wave no longer propagates parallel to $\vec{Q}$, but parallel to $\vec{E}_i$.[11] Thus not only would ν_{an} increase, but also there would be no need to make any approximations concerning the angular width of the spectrum. In a future publication, we plan to discuss the problem of absorption in a magnetic field.

Equations (1)–(4) are a coupled set of equations which we solve numerically starting at $x=0$ and integrating backwards towards the laser. We assume that the density profile near the critical density has steepened[8,12] until reaching the steady state described in Ref. 12. We start with $\sum|e\varphi/T_e|^2=10^{-5}$. At the initial location ($x=0$) in the underdense plasma the flow velocity v, and V_{0s}/V_e [$\equiv e(E_i+E_r)/m\Omega V_e$] are determined in terms of the density.[12] Also $E_i^2=E_r^2$ at $x=0$. Choosing an electron temperature is then essentially equivalent to choosing an incident laser power. The remaining initial parameters to be specified are T_i/T_e and Q. The parameter $Q(x=0)$ is found by iteration so that $Q(x=-\infty)=0$.

Results are shown in Figs. 1(a)-1(d). Figure 1(a) shows the spatial dependence of T_e, Q, E_i^2, E_r^2, and $[\sum|e\varphi(k)/T_e|^2]^{1/2}$, where $T_e(x=0)=12$ keV, $T_i(0)/T_e(0)=\frac{1}{30}$, $\omega_{pe}^2(0)/\Omega^2=0.7$, and $\theta=0$. Following E_i^2 back to $-\infty$, we see that the incident laser flux is 10^{15} W/cm^2, $Q(0)/nmv_e^3(0)\approx 0.1$, and about $k\sim k_D/2$ and $e\varphi/T$ maximizes at a value of 0.2, which seems just barely small enough that one can believe linear theory. The electrons which principally absorb the laser light have velocity $\sim\Omega/k\sim 3v_e$ so that an energetic tail is not expected to be substantially produced.[13] Figure 1(b) shows the absorption efficiency as a function of density at $x=0$ assuming $T_e(0)=12$ keV and $T_e(0)/T(0)=30$. Figure 1(c) shows values of $Q(0)/nmv^3(0)$, and absorption efficiency as a function of laser power where $n(0)=0.5n_{cr}$ (n_{cr} is the critical density) and $T_e(0)/T_i(0)=30$. The high temperatures calculated here (and high temperature ratios assumed here) of course, exist only in front of the critical surface. At higher densities, the temperature would be much lower. Furthermore two temperature fluid codes show high temperature ratios in the underdense plasma.[14] Figure 1(d) shows the absorption efficiency as a function of angle for $n(0)=0.5n_{cr}\cos^2\theta$ (see Ref. 12), $T_e(0)/T_i(0)=30$ and $T_e(0)=12$ keV (the incident laser flux was in the vicinity of 2×10^{15} W/cm^2). The fractional absorption would be substantially increased at higher power and/or with higher density at $x=0$ as is apparent from Figs. 1(b) and 1(c). Notice also that the inward energy flux due to thermal conduction is comparable to the outward thermal energy flux due to hydrodynamic expansion.

A potential problem is that the spatial scale length shown in Fig. 1(a) may be too long. This long length allows for Brillouin scattering. Also it is not clear that the plasma can expand this far in a short-pulse experiment. We are currently investigating means by which this length can be reduced.

To summarize, our results show good absorp-

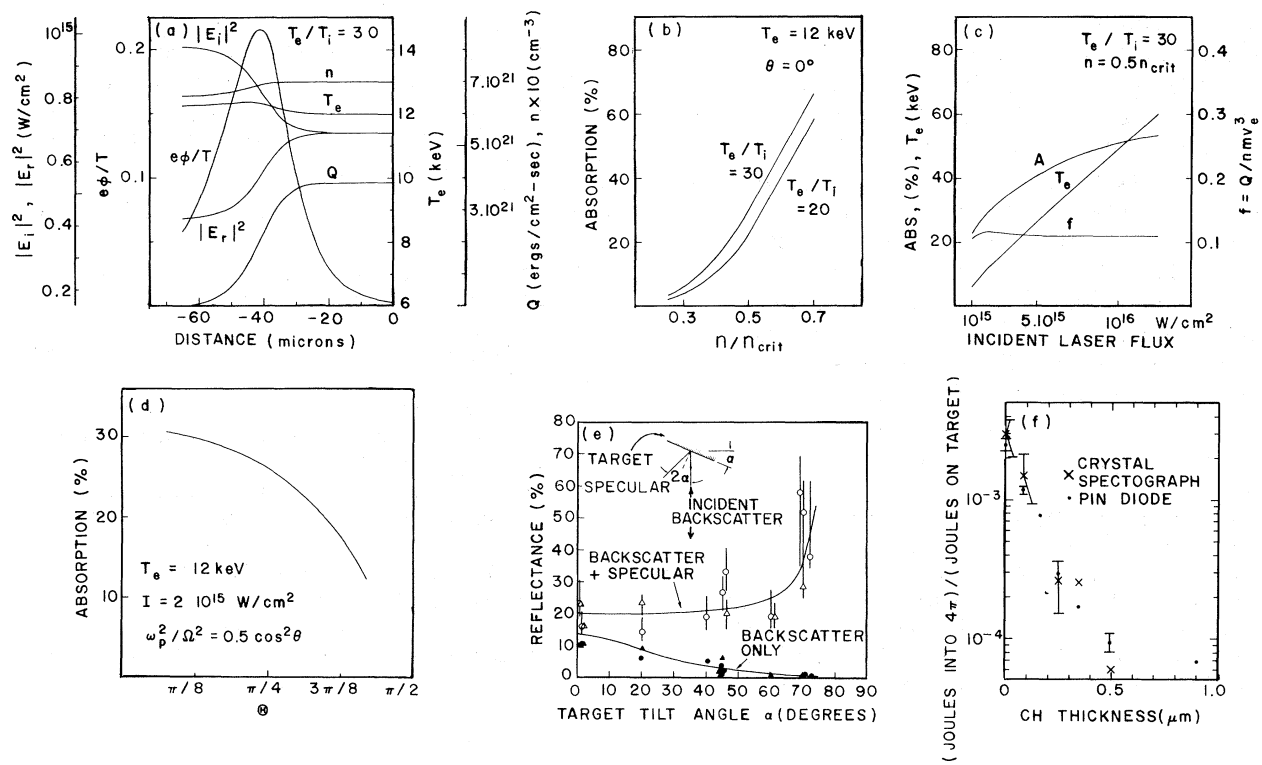

FIG. 1. (a) Fluid quantities as a function of x; (b) fractional absorption as function of density; (c) fractional absorption, T_e and f as a function of laser flux; (d) fractional absorption as a function of tilt angle; (e) total reflected and backscattered energy as a function of the angle of incidence (circles and triangles are for the two incident polarizations); (f) Al line radiation (×) and contiuum (·) (1–3 keV) vs polystyrene thickness.

tion by the thermal part of the distribution function which is nearly independent of both incident polarization and angle, and with Q/nmv_e^3 of typically about 0.1. Finally, we wish to point out that similar results were found by others.[15]

We will now discuss more fully some of the related experimental results. Figure 1(e) (taken from Ref. 5) shows the reflected light as a function of angle of incidence and polarization, for laser irradiance of 5×10^{15} W/cm^2. Notice that the absorption efficiency is not strongly dependent on either polarization or angle of incidence for tilt angle less than about 60°, just as predicted in Fig. 1(d).

We now discuss the relevance of an experiment in which the transport of energy was studied at an irradiance of 10^{15} W/cm^2 through a thin layer of polystyrene into an aluminum substrate.[4] The intensity of aluminum line radiation is shown as a function of polystyrene thickness in Fig. 1(f). The absorbed energy flux Q is measured. Assuming a particular value for Q/nmv_e^3 then allows one to calculate the average electron energy near the critical surface. Since the laser energy must ultimately heat the electrons in the polystyrene to this energy, one can calculate how much laser energy, at given irradiance, is needed to just burn through a given layer of polystyrene (i.e., to cut off the aluminum line radiation).

A value of Q/nmv_e^3 of about 0.1 is consistent with the upper limit of $Q/nmv_e^3 \sim 0.2$ inferred from the dependence of Al x-ray radiation to polystyrene thickness shown in Fig. 1(f).[4]

The experimental situation is closely one-dimensional. The asymmetry of specularly reflected light indicates that, on the average, the center of the critical surface bulges by only about 1 μm compared to the half-energy-content focal diameter of 30 μm. However, Fourier analysis of the specularly reflected light indicates possible enhanced density fluctuation near 1 μm,[7] which is close to the peak-ion-fluctuation wavelength.

Finally, we would like to make a few remarks on hard-x-ray data. Our own[16] and other theories[8] have shown that resonant absorption creates electron distributions having nonthermal tails extending from about $3v_e$ to 6 or $7v_e$. If 10^{15} W/cm^2 is conducted by the electrons and $Q/nmv_e^3 \sim 0.2$ as indicated in Ref. 4, then the temperature is about 6 keV. Thus the nonthermal tail would extend from about 60 keV to about 300 keV. The layered-target experiments and others[6] show very few hard x rays above 100 keV. Thus, there appears to be no indication of a strong superthermal tail to the electron distribution function.

In summary then, there are good theoretical and experimental indications that light absorption by enhanced ion-density fluctuations is a very important process for laser fusion.

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Phase Transition on Mo(100) and W(100) Surfaces

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Low-energy-electron-diffraction studies of carefully cleaned and annealed surfaces of molybdenum (100) and tungsten (100) show that a phase transition can be induced by lowering the temperature below 300 K. The periodicity of the "reconstructed" surface is believed to be due to the formation of a displacement wave with a wavelength which is $2a$ (a is the lattice parameter) for W(100) and $\sim 2.2a$ for Mo(100). The phase transition is reversible and seemingly second order. It appears possible that displacements of this type also occur in chemisorption on these surfaces.

In recent low-energy-electron-diffraction (LEED) studies we have observed a temperature-dependent structural transformation occurring on the molybdenum (100) surface. The transformation is completely reversible and appears to be characteristic of the clean surface. This discovery has led us to re-examine the behavior of the tungsten (100) surface which is one of the most widely studied substrates and which in many respects is similar to Mo(100). As discussed below, the results indicate that a "reconstruction" of the clean surface takes place also in the case of W(100) as the temperature is lowered.

The relevant LEED patterns are shown in Fig. 1. Above room temperature the pattern from either Mo(100) or W(100) shows only the "normal" spots [Fig. 1(a)], but as the crystal is cooled in vacuum the pattern changes. On Mo(100) the change consists of the appearance of a quartet of spots around the $(\frac{1}{2}\frac{1}{2})$ positions [Fig. 1(b)]; on W(100) single $(\frac{1}{2}\frac{1}{2})$ spots appear [Fig. 1(c)]. The variation of the intensity of the extra diffraction spots with temperature is shown in Fig. 2. As seen, the intensity changes gradually, suggesting a second-order transition. No hysteresis was seen for either sample as the temperature was raised or lowered. For both surfaces, as the intensity decreases with increasing temperature, the extra diffraction spots become larger, more diffuse, and streaky, presumably because of critical scattering. On the basis of these observations and data obtained by other techniques we