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Thermal heat-flux reduction in laser-produced plasmas

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It is shown that the non-Maxwellian heating by inverse bremsstrahlung which results in a truncation of the tail of the electron-distribution function is responsible for a reduction of the thermal conductivity in a smooth temperature gradient by a factor of 3 to 4 from the Spitzer-Härm conductivity. In a steep temperature gradient, the reduction is still a factor of about 1.5 from the recent nonlinear theories based on a Maxwellian electron-distribution source.

Thermal conduction plays an important role in the transport of energy in laser-heated targets. Interpretation of many experimental results suggests that the heat flux may be much smaller than expected on the basis of the classical Spitzer-Härm[1] (SH) theory.[2] Actually it is currently admitted that the heat flux cannot exceed a small fraction (typically 0.03) of the so-called free streaming value $q_f = n_e m_e v_e^3$, where $v_e = (T_e/m_e)^{1/2}$ is the thermal velocity. One of the most convincing explanations is based on the study of the validity of the SH theory in steep temperature gradient. The SH theory is based on a Chapman-Enskog development[3] which assumes that the typical collision mean free path λ is smaller than the typical temperature scale length L_t. It has been pointed out by Gray and Kilkenny[4] that in a completely ionized gas, the heat flux is carried mainly by particles whose velocity is between 3 to 4 times the thermal velocity v_e, and whose electron-ion collision mean free path $\lambda(v) = \lambda_e(v/v_e)^4$, is 100 to 300 times the thermal mean free path λ_e. As a result, the SH theory is limited to very small temperature gradients, $\lambda_e/L_t \lesssim 2.10^{-3}$. Numerical simulations[5–7] of heat transport in laser-produced plasma indicates that in a steep temperature gradient, the thermal heat flux on the main body of the heat front is reduced by roughly one order of magnitude below that given by the SH description and is limited to about $0.1\,q_f$ while the thermal conductivity at the base of the heat front exceeds the SH conductivity, because of the nonlocal heat transport due to hot, nearly collisionless, electrons streaming away from the top of the heat front.

However, the SH theory and the numerical simulations of Refs. 5–7 assume that the source-electron-distribution function is a Maxwellian distribution. If the light absorption is due to inverse bremsstrahlung, Langdon[8] has demonstrated that the electron-distribution function is far from a Maxwellian distribution if $Z v_{\rm osc}^2/v_e^2 \gtrsim 1$, where $v_{\rm osc}$ is the peak velocity of oscillation of the electrons in the high-frequency electric field, and Z is the ionization state. The resulting "top-hatted" distribution function is depopulated in the velocity range $v > 3v_e$, which determines the heat conductivity in the SH theory. We demonstrate in this Brief Report that such non-Maxwellian effects may result in a reduction of the thermal conductivity by a factor of 3 to 4 in the linear regime. This result is in agreement with the work of Dum[9] about the anomalous transport due to ion sound (the electron heating due to an isotropic turbulent ion sound spectra and to inverse bremsstrahlung leads to the same electron-distribution function[8,10]). We also demonstrate that in a steep temperature gradient, the heat transport is still reduced compared with the Maxwellian source case.

We first suggest that this reduction is significant even at low intensities ($Z v_{\rm osc}^2/v_e^2 \gtrsim 0.1$) because electron-electron collisions are not rapid enough to fill out the high-velocity tail of the Maxwellian distribution, which is responsible for the heat transport. As a clue to this result, consider the ratio of the e folding time for heating of the bulk plasma to the e-e equilibration time τ_{ee} required to establish a Maxwellian distribution in the range $(3-4)v_e$. The heating time is shorter than τ_{ee} when $Z v_{\rm osc}^2/v_e^2 \gtrsim 0.1$, in which case non-Maxwellian heat transport is possible. Note this limit is considerably lower than in the absorption problem,[8] where the absorption is due to the bulk of the distribution function, mainly because of the strong scaling of τ_{ee} with the velocity ($\tau_{ee} \propto v^3$). Furthermore, the heat transport itself tends to depopulate the hot part of the distribution function of the corona of laser-produced plasma.

To calculate the heat transport properties of a non-Maxwellian plasma, we use a test isotropic-distribution function of the form[9]

$$f_0(v) = [n/4\pi\Gamma(3/n)](n_e/v_0^3)\exp[-(v/v_0)^n] \quad . \quad (1)$$

$n=2$ for a Maxwellian distribution function, $n=5$ for a laser-heated-distribution function when e-e collisions are neglected,[8] and for a plasma heated by ion-sound turbulence,[10] and $n=\infty$ corresponds to the water bag model. The temperature T_e is given by

$$T_e = m_e v_e^2 = m_e \langle v^2/3 \rangle = [\Gamma(5/n)/3\Gamma(3/n)] m_e v_0^2 \ .$$

The total electron-distribution function is expanded as $f = f_0 + \vec{f}_1 \cdot \vec{v}/v$.[11] In the presence of small gradients, the kinetic equation for f_1 is given by

$$\vec{f}_1 = -\lambda(v)[\nabla f_0 - (e/m_e v)(\partial f_0/\partial v)\vec{E}] \ . \qquad (2)$$

Using the zero-current condition, one obtains

$$\frac{f_1}{f_0} = \left[\frac{n}{2}u^n - \frac{5n}{12}\frac{\Gamma(8/n)}{\Gamma(6/n)}u^{n-2} - \frac{3}{2}\right]\frac{\lambda(v)}{L_t} - \left[\frac{n}{6}\frac{\Gamma(8/n)}{\Gamma(6/n)}u^{n-2} - 1\right]\frac{\lambda(v)}{L_n} \ , \qquad (3)$$

where $u = v/v_0$, $L_t = T/(dT/dx)$, and $L_n = n_e/(dn_e/dx)$. For Coulomb scattering $\lambda(v) = \lambda_e(v/v_e)^4$ and the heat flux q is given by

$$q/q_f = -(K_t\lambda_e/L_t + K_n\lambda_e/L_n) \ , \qquad (4)$$

where

$$K_t = a(7b - 5c) \ , \quad K_n = 2a(b-c) \ ,$$

$a = [\Gamma(3/n)]^{5/2}[3/\Gamma(5/n)]^{7/2}$, $b = \Gamma(10/n)/12$, and $c = [\Gamma(8/n)]^2/9\Gamma(6/n)$. These transport coefficients agree with the ones given by Dum.[9] As already pointed out by Dum[9] and Shkarofsky[11] we note that there is a heat-flow term associated with the density gradient if $n \neq 2$. This is linked with the fact that non-Maxwellian distribution functions ($n \neq 2$) of the form (1) do not correspond to steady-state solutions of the Vlasov equation when $\nabla T_e = 0$ and $\nabla n_e \neq 0$. In Fig. 1(a), we have plotted the transport coefficient as functions of n. We observe that the linear thermal conductivity is divided by about 4 between $n=2$ (Maxwellian) and $n=5$ (inverse bremsstrahlung heating). If one assumes that the pressure is almost constant on the body of the heat front, $L_n \simeq -L_t$, and the heat flux is divided by about 3 between $n=2$ and 5.

We now evaluate the limit of validity of the linear theory for a finite temperature gradient. We expect the linear theory to be valid if $|f_1/f_0| < 1$ for the particles which effectively carry the heat flux. To obtain a precise criterium, we ask $|f_1/f_0| < 1$ for the velocity v^* corresponding to the maximum of the differential heat flux $Q(v) = (4\pi m_e/6)v^5 f_1\, dv$, where $q = \int_0^\infty Q(v)dv$. This velocity is such that $Q(v^*) > 0$ and $\partial Q(v^*)/\partial v = 0$. [See the inset of Fig. 1(b).]

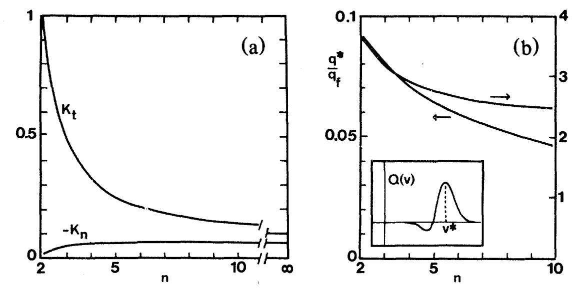

FIG. 1. (a) Transport coefficients $K_t/K_t^{(2)}$ and $K_n/K_t^{(2)}$ as functions of the exponent n in Eq. (1). $K_t^{(2)}$ is the usual Spitzer-Härm value of K_t, corresponding to $n=2$. (b) Characteristic velocity v^* corresponding to the maximum of the differential heat flux $Q(v)$, and characteristic heat flux q^* obtained when the temperature gradient is such that $|f_1/f_0| = 1$ for $v = v^*$. $q_f = n_e m_e v_e^3$ is the free streaming limit. n is the exponent in Eq. (1).

The value of v^*/v_e when $\nabla n_e = 0$ is reported in Fig. 1(b) as a function of n. We note that, as n increases, the heat flux is due to electrons of comparatively lower velocities: $v^*/v_e = 3.71$ for $n=2$ and $v^*/v_e = 2.75$ for $n=5$. This suggests that the linear theory is valid for comparatively larger temperature gradients. In fact, computing λ_e/L_t such that $|f_1/f_0| = 1$ for $v = v^*$, we obtain $(\lambda_e/L_t)^* \simeq 0.002$ for $n=2$ and $(\lambda_e/L_t)^* \simeq 0.005$ for $n=5$. The corresponding heat flux are $q^*/q_f = 0.093$ for $n=2$ and $q^*/q_f = 0.064$ for $n=5$. q^* is reported in Fig. 1(b) as a function of n.

When $\lambda_e/L_t \gtrsim (\lambda_e/L_t)^*$ a nonlinear and nonlocal theory is necessary.[5–7,12] We expect that the thermal heat flux on the main body of the heat front will be a function of the ratio λ_e/L_t, as suggested by Shvarts *et al.*[6] in the $n=2$ case. A possible approximation is the harmonic mean function q_h given by $q_h^{-1} = (q^*)^{-1} + q^{-1}$, where q is given by Eq. (4). On the other hand, the heat flux at the base of the heat front, though exceeding the linear heat flux given by Eq. (4), because of the nonlocal heat transport due to the hot electrons streaming away from the top of the heat front, is smaller in the $n=5$ than in the $n=2$ case, because of the depletion of the distribution function in the high-velocity range.

In summary, the non-Maxwellian distribution from inverse bremsstrahlung results in a factor of 1.5 (large temperature gradients) to 4 (small temperature gradients) reduction in transport, which may explain submicron wavelength laser experiments.[13]

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