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1.E Two-Dimensional, Nonlocal Electron Transport in Laser-Produced Plasmas

Much effort has been devoted to the study of nonlocal electron transport in laser-produced plasmas.[1,2] Most of the work has involved the numerical solution of the electron Fokker-Planck (FP) equation in one dimension by finite difference techniques. Comparisons with standard fluid transport calculations using Spitzer-Härm (SH)[3] heat flow $\underset{\sim}{q}_s = -\kappa \underset{\sim}{\nabla} T$ has revealed the occurrence of hot-electron "preheat," as well as reduced penetration of the bulk heat front, a phenomenon known as "flux inhibition."[1] To simulate the latter effect, fluid codes are normally equipped with an artificial flux limiter, which maintains the heat flux $\underset{\sim}{q}$ below some fraction f of its free-streaming limit $q_f = nm(T/m)^{3/2}$, i.e., $\underset{\sim}{q} = \underset{\sim}{q}_s/(1+|\underset{\sim}{q}_s/fq_f|)$.[4,5]

This technique has been shown to be adequate for long-pulse ($\sim$1-ns), short-wavelength ($<$1-μm) lasers at moderate irradiances ($<10^{15}$ W/cm^2), where usually $q_s \lesssim fq_f$ (for $f = 0.1$–0.2).[6] However, if the transport is two-dimensional (2-D), as would be the case for nonuniformly illuminated laser plasmas, the validity of fluid theory has not yet been tested. In particular, the application of flux limiters requires special care, since the direction of the heat flow may not always be parallel to $-\nabla T$.[5,7] Nevertheless, 2-D fluid codes are widely used for thermal transport studies with each direction individually flux limited.[8]

Here, we investigate the validity of the fluid modeling by constructing a code (SPARK) designed to numerically solve the electron FP equation in 2-D planar geometry under conditions relevant to laser-produced plasmas. By imposing a spatial inhomogeneity of scale ℓ in the incident laser beam, we show that the nonlocal electron transport is less effective at smoothing temperature gradients than the corresponding diffusive transport based on SH heat flow, irrespective of the size of the flux limiter. This effect becomes significant when $\ell \lesssim 80\lambda_{mfp}$, where λ_{mfp} $(= 3T^2/4\sqrt{2\pi}\, Zne^4 \ln\Lambda)$ is the electron-ion mean free path at the critical surface.

The numerical algorithm of SPARK is briefly described as follows (for further details see Ref. 7). Using the diffusive approximation and taking the high-Z limit of the FP equation, we obtain (using the notation of Shkarofsky, Johnston, and Bachynsky)[9]

$$\frac{\partial f_0}{\partial t} + \frac{v}{3}\nabla\cdot \underset{\sim}{f}_1 = \frac{1}{v^2}\frac{\partial}{\partial v}\left[\frac{v^2}{3}\,\underset{\sim}{a}\cdot \underset{\sim}{f}_1 + Y\left(Cf_0 + D\frac{\partial}{\partial v}f_0\right) + \frac{YnZ}{6v}v_0^2\frac{\partial}{\partial v}f_0\right], \tag{1a}$$

$$\underset{\sim}{f}_1 = -\tau\left(v\nabla f_0 - \underset{\sim}{a}\frac{\partial}{\partial v}f_0\right), \tag{1b}$$

where

$$\underset{\sim}{a} = \frac{|e|}{m}\underset{\sim}{E},\quad \tau = v^3/[(Z+1)\,nY],\quad Y + 4\pi\,(e^2/m)^2\ln\Lambda,\quad C = I_0^0 f_0,$$

$$\text{and } D = \frac{v}{3}\left(I_2^0 + J_{-1}^0\right).$$

To simplify the analysis, we have assumed the ions to be cold and motionless. The last term in Eq. (1a) is the inverse-bremsstrahlung operator, and v_0 is the electron quiver velocity.[10] This reduced form of the FP equation has been successfully used for 1-D transport studies,[2] where one is not too interested in highly collisionless phenomena (such as Landau damping). Here, we have further neglected the effect of magnetic fields that may be justified by the fact that we will be comparing our results with fluid calculations based on the same approximation. Furthermore, we will simulate the transport under conditions where the effect of magnetic fields on the transport is not expected to be significant, i.e., short laser wavelength at moderate irradiances.[5]

Substituting Eq. (1b) into Eq. (1a) and defining

$$\alpha = -[\partial_v \ln(f_o)]/v,\ \underset{\sim}{\beta} = -\underset{\sim}{\nabla}\ln(f_o) \text{ and } \chi = v^2\tau/3,$$

we have

$$\begin{aligned}\frac{\partial f_o}{\partial t} &= \underset{\sim}{\nabla} \cdot [\chi\,(\underset{\sim}{\nabla} f_o + \underset{\sim}{a}\alpha f_o)] \\ &+ \frac{1}{v^2}\frac{\partial}{\partial v}\left\{\chi\left(\underset{\sim}{a}\cdot\underset{\sim}{a}\frac{\partial}{\partial v} f_o + v\underset{\sim}{\beta}\cdot\underset{\sim}{a}\, f_o\right)\right. \\ &+ \left. Y\left[Cf_o + \left(D + \frac{nZv_o^2}{6v}\right)\frac{\partial}{\partial v} f_o\right]\right\}. \end{aligned} \qquad (2)$$

Assuming that the nonlinear coefficients α and $\underset{\sim}{\beta}$ are known explicitly, as is usually done with the terms $\underset{\sim}{a}$, C, and D, we have transformed the troublesome cross-derivatives into convective terms. The resulting diffusion-convection equation may now be solved by standard numerical techniques.

Equation (2) is differenced in conservative form by adopting Cartesian geometry in x, z, and v. A generalized Chang-Cooper[11] weighting is applied in all directions by using the zero flux condition across the cell boundaries. The electric field is calculated via the "implicit moment"[12] method by assuming that $\underset{\sim}{\nabla} \times \underset{\sim}{E} = 0$ and taking the

$$\int_0^\infty v^2 dv$$

moment of Eq. (2). (An alternative scheme that assumes the total current is zero and allows for a finite $\underset{\sim}{\nabla} \times \underset{\sim}{E}$ has also been tried with little effect on the transport calculations, apart from a deterioration in the quasi-neutrality.)[13] The resulting system of equations is solved by an "alternating-direction-implicit"[7,14] scheme in v, z, and x respectively, in conjunction with a "predictor-corrector" step. In the predictor stage, f_o is linearly extrapolated to the half-time level and used to calculate the nonlinear coefficients. In the corrector stage, f_o is averaged in time to recalculate the coefficients.

To illustrate the process of 2-D transport we consider a planar, fully ionized plasma of $Z = 4$ with an initial temperature of 250 eV. The ion background is kept fixed throughout the simulation with an exponential ramp of 25 μm, as shown in Fig. 36.21. A 0.35-μm laser is propagated along the positive z direction and is absorbed via a 1-D ray-trace package, as used in the hydrocode *LILAC*,[15] with a full reflection at the critical surface. The intensity of the beam is modulated in the x direction by $I = I_o[1 + \epsilon\cos(2\pi x/\lambda)]$ where $I_o = 5 \times 10^{14}$ W/cm^2 at all times.

Figure 36.21 shows the density, temperature, and absorption profiles of a 1-D simulation at 120 ps, using SPARK and a fluid code with (a) $f = \infty$ (i.e., no flux limit) and (b) $f = 0.2$. As expected, for UV light at moderate irradiances, the agreement between *FP* and *SH* calculations is very good, and the mild flux inhibition is fairly well modeled by $f = 0.2$, in agreement with Ref. 6.

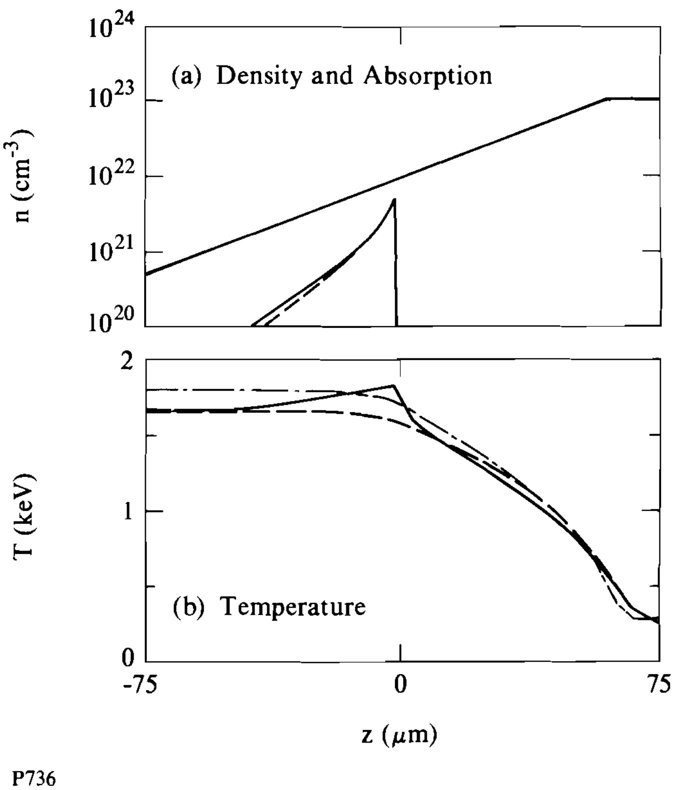

Fig. 36.21
(a) Density and energy deposition (in relative units), (b) temperature profiles for 0.35-μm irradiation at 5×10^{14} W/cm^2 after 120 ps. Solid curves refer to *FP* results; dashed and dash-dotted curves refer to *SH* results with $f = \infty$ and $f = 0.2$, respectively.

The result of nonuniform illumination is shown in Fig. 36.22 for $\epsilon = 1$ and $\lambda = 150$ μm, where the comparison is made with fluid results under no flux limitation. Here, we used (30 $\times$ 30) cells in the x-z plane and obtained a maximum charge separation of 5×10^{-5} and an energy conservation error of 6×10^{-3} for a time step of 0.1 ps. As observed from the isotherms, the smoothing predicted by the *FP* solution is less than predicted by the *SH* solution, though the average temperature $<T>$ follows the 1-D result of Fig. 36.21 very closely. This effect is highlighted in Fig. 36.23, where we plot

$$\sigma_{\mathrm{rms}} = \{ [\int dx (T - <T>)^2] / \int dx \}^{1/2} / <T>$$

as a function of z. The peaked structure of σ_{rms} at high density is essentially an artifact of the locally sharp temperature gradients in z. Paradoxically, the simulation predicts $|q_x| < |q_z|$ and $|q_x| << fq_{fs}$ (apart from the very-low-density corona), which is normally taken to

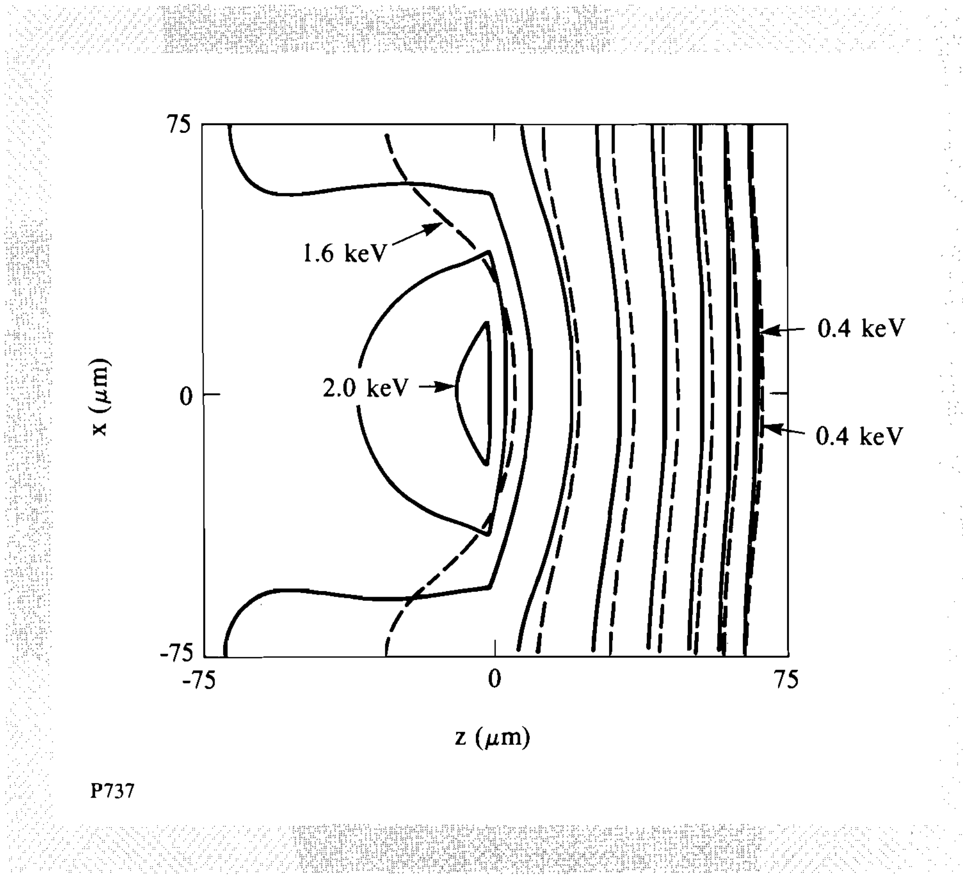

Fig. 36.22
Isotherms for nonuniform laser irradiation in steps of 0.2 keV (as in Fig. 36.21).

imply that the transverse heat transfer is close to classical. However, as has been previously shown by Bell[16] in the context of 1-D thermal transport in ion waves, the relevant criterion for nonlocal effects to be significant is that the mean free path of the heat-carrying electrons ($\approx 80\lambda_{mfp}$) is greater than the relevant spatial scale length ℓ ($= \lambda/2\pi$), irrespective of the magnitude of the heat flux. This is demonstrated quantitatively in Fig. 36.24 where we plot the ratio of *FP* to *SH* $\sigma_{\rm rms}$ as a function of ℓ/λ_{mfp}, calculated at the critical surface. The observed reduction in smoothing by the nonlocal transport is a direct consequence of heat-flux inhibition and is independent of ϵ (for $\epsilon \leq 1$). Similar results have also been obtained for different laser wavelengths and irradiances.[13] It must be noted, however, that although $(\sigma_{\rm rms})_{FP} >> (\sigma_{\rm rms})_{SH}$ for $\ell << \lambda_{mfp}$, both models predict that $\sigma_{\rm rms} \rightarrow 0$ in the same limit, as seen on Fig. 36.24.

The effect of flux limiting the *SH* heat flow in 2-D is shown in Fig. 36.23 for $f = 0.2$. The relatively small improvement in the results is not surprising since the flux limitation is only weakly dependent on the lateral heat flow, which is well below fq_f. Moreover, the increased coronal temperature arising from the axial flux inhibition (see Fig. 36.21) has a tendency to increase the thermal conductivity coefficient and, hence, the amount of smoothing.

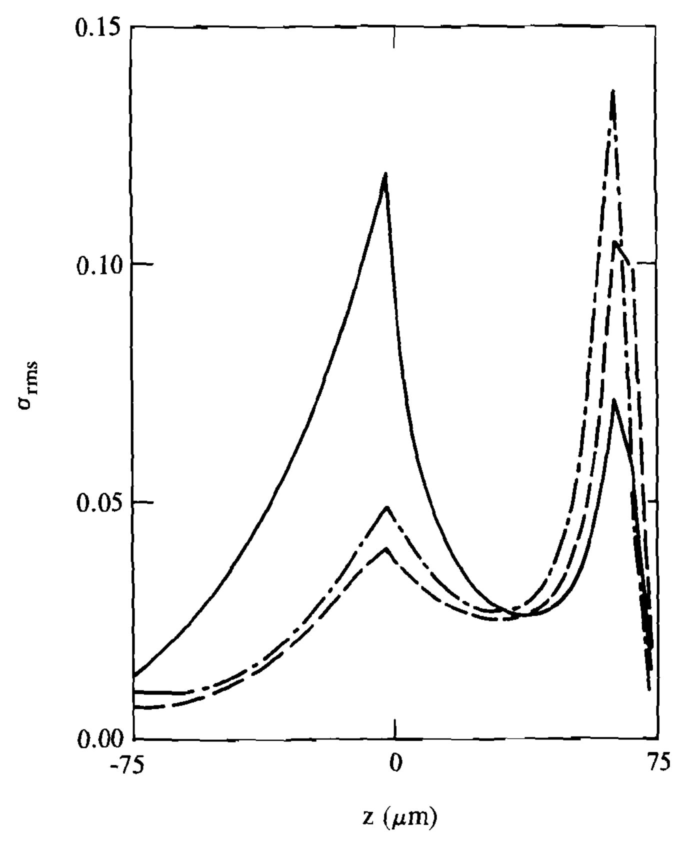

Fig. 36.23
Plot of σ_{rms} as a function of z (as in Fig. 36.21).

In summary we have shown that even when 1-D transport in laser plasmas is well described by fluid equations the same is not necessarily true in 2-D. Particularly, when there are small-scale ($\lesssim 80\lambda_{mfp}$) temperature modulations in the plasma, the thermal smoothing becomes less effective as a result of the nonlocal nature of the electron transport. This reduced smoothing, which is not adequately treated by a simple flux-limited diffusion theory, may have important implications to the process of thermal self-focusing and ultimately to the estimation of ablation pressure uniformity. In order to adequately assess these effects, SPARK has been recently extended to include the hydrodynamic response of the plasma, as well as the refraction and diffraction of the laser beam in the corona.

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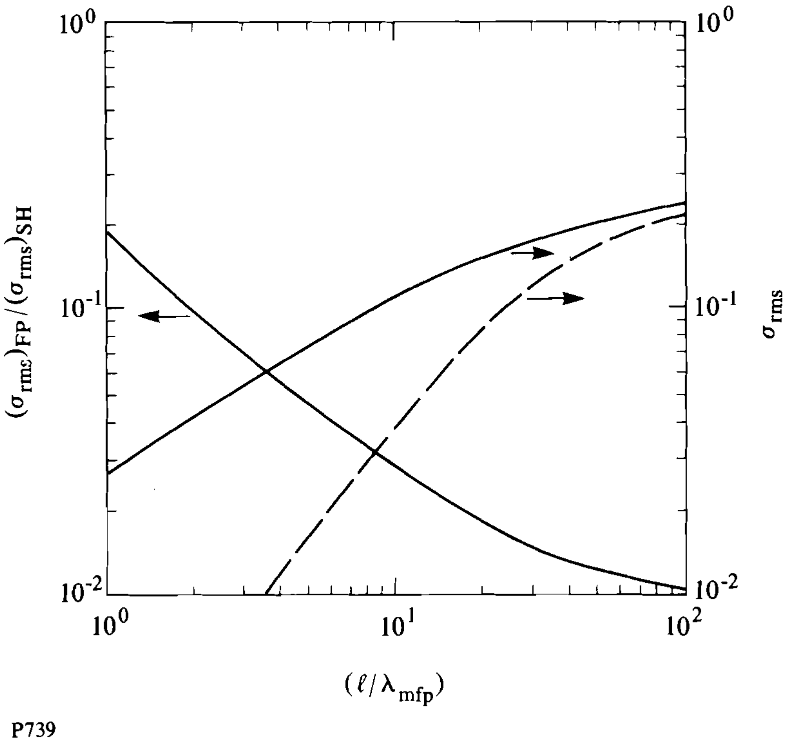

Fig. 36.24
Plots of $(\sigma_{rms})_{FP}$, $(\sigma_{rms})_{SH}$, and $(\sigma_{rms})_{FP}/(\sigma_{rms})_{SH}$ as functions of ℓ/λ_{mfp}, calculated at the critical surface (as in Fig. 36.21).

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