sults of counter-injection experiments in hydrogen support these conclusions. They also show that the argon radiation from the interior increases, although the maximum emission is a factor of 3–5 lower than in deuterium.

Evaluation of the co-injection results is not so straightforward, because either reductions of the accumulative mechanisms or enhancement of anomalous diffusion can prevent the accumulation. But the fact that neoclassical, beam-induced effects seem to increase the accumulation rates of impurities during counter-injection, even in hydrogen discharges where the anomalous processes are relatively large, provides a strong argument that the same mechanisms inhibit the impurity influx during co-injection. In addition, the magnetohydrodynamic activity, which often serves as an indicator of changes in the anomalous impurity transport, remains at a very low level. It is not significantly intensified by injection. Therefore, it appears that increases of anomalous transport are of secondary importance in the experiments we describe here, and the lack of accumulation during co-injection is caused by reductions of the neoclassical effects.

In summary, the present results from the ISX-*B* tokamak indicate that co-injection inhibits inward impurity transport and counter-injection enhances it, in qualitative agreement with recent theories.[1,2] The capability of making these observations requires that anomalous diffusive effects do not overwhelm neoclassical mechanisms and that increases of impurity influxes during injection do not obscure transport changes.

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Apparent and Real Thermal Inhibition in Laser-Produced Plasmas

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Laser-driven thermal-electron-transport inhibition is studied with use of a self-consistent Monte Carlo model for all the electrons. Comparison is made with the results of flux-limited single-group–diffusion calculations. The need for severe flux limiters is traced to deficiencies in the classical diffusion modeling that can excessively heat the overdense surface matter of a pellet, ignore coronal decoupling of the thermals electrons, and neglect the effects of electric-fields for a return current through density gradients.

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For some time it has been recognized[1,2] that severe flux limitation must be imposed on thermal electron diffusion for accurate modeling of laser-plasma interactions. The implied thermal inhibition has been ascribed to ion-acoustic turbulence,[3] suprathermal and thermal counter-streaming,[4,5] and transition-collisional effects.[6] Yet in simulation, the ion-acoustic turbulence is too weak[7,8]; in practice, the inhibition has been evident even at low intensities where the counter-streaming is minimal; and the transition-regime calculations, although promising, have yielded a flux limit factor of $f = 0.1$—which is, at least, three times too large.

These considerations have encouraged a further examination of the transition regime. In this Letter for the first time we compare the transport predicted by flux-limited single-group thermal diffusion to transport from a fully self-consistent Monte Carlo model. We show that the classically limited diffusion ($f = 0.6$) heats the overdense surface mass of laser targets to too high a temperature, especially in the presence of a steep density rise near critical, while the Monte Carlo transport heats much more mass in depth to significantly lower temperatures. The latter should result in constrained hydrodynamics and reduced radiative emissions, in agreement with experiment. Further, we show that with $f \simeq 0.03$ the limited-diffusion solutions are brought into some accord with the Monte Carlo results. Thus, the need for strong thermal flux limitation is identified as an artifact of our efforts to employ single-group diffusion to model transport that is intrinsically multigroup.

Many of the features of our earlier hybrid scheme[5,8] have been retained, except that now the ions and *all* the electrons are weighted particles. The weights are adjusted to match any initial density profile. The electrons are heated by giving them velocity increments from a Gaussian distribution at the heater temperature, $k\Delta T = \frac{2}{3}\mathcal{E}$, where $\mathcal{E}$ is the absorbed particle energy per cycle cell. Heating is terminated when some prescribed temperature is achieved, or, alternatively, $\mathcal{E}$ is established in consistency with inverse-bremsstrahlung absorption up to the critical density. The electrons are scattered with the mean-squre deflection angle $\langle\theta^2\rangle = 8\pi e^4 m^{-2} c^{-3}\Delta t Z n_i (Z+1)\ln\lambda \equiv 2\nu(c)\Delta t$, as in Refs. 5 and 8. Here $c^2 = u^2 + v^2$ and v is the transverse velocity. However, now each electron loses energy, and has its speed reduced by Coulomb drag against the other electrons, in accordance with $\Delta c = -4\pi e^4 (mkT_e)^{-1} \times \Delta t\, G(\xi)\ln\lambda$, $\xi = (mc^2/2kT_e)^{1/2}$, where $G(\xi) = 0.376\xi \times (1 + 0.542\xi - 0.504\xi^2 + 0.752\xi^3)^{-1}$ is a polynomial fit[9] to Spitzer's error function combination. The lost energy is redeposited isotropically in the drift frame of the electrons by the usual heating procedure. The E field is now calculated by the implicit moment method,[10] which gives greater accuracy than plasma period dilation.[5,8] For this, following the heating, scatter, and drag of particles, we accumulate the fluid density, flux, and total pressure moments $n_\alpha^{(m)}$, $j_\alpha^{(m)}$, and $P_\alpha^{(m)}$ ($\alpha = h, c$, and i for hot and cold electrons and ions, respectively). Then, we use the momentum, continuity, and Poisson equations to predict values ($\tilde{x}$) for the moments and fields at the end of the next cycle. Thus, storing all the information at level (m), we obtain

$$\tilde{j}_\alpha^{(m+1/2)} = j_\alpha^{(m)} - \frac{1}{m_\alpha}\left(\frac{\partial P_\alpha^{(m)}}{\partial x} - q_\alpha n_\alpha^{(m)} E^{(*)}\right)\frac{\Delta t}{2}, \tag{1a}$$

$$\tilde{n}_\alpha^{(m+1)} = n_\alpha^{(m)} - \frac{\partial}{\partial x}(\tilde{j}^{(m+1/2)})\Delta t, \tag{1b}$$

$$\tilde{E}^{(m+1)} = 4\pi\int_0^x \sum_\alpha q_\alpha \tilde{n}_\alpha^{(m+1)}\,dx + \tilde{E}^{(m+1)}(0), \tag{1c}$$

which, with the centering $E^{(*)} = (\tilde{E}^{(m+1)} + \tilde{E}^{(m)})/2$, and the assumption $E(0) = j(0) = 0$ for a quiescent left boundary, rearranges to

$$\tilde{E}^{(m+1)} = 4\pi\left\{\int_0^x \sum q_\alpha n_\alpha^{(m)}\,dx - \sum q_\alpha j_\alpha^{(m)}\Delta t + \sum\frac{q_\alpha}{m_\alpha}\frac{\partial P_\alpha^{(m)}}{\partial x}\frac{(\Delta t)^2}{2} - \frac{\omega_p'^2(\Delta t)^2}{4}\tilde{E}^{(m)}\right\}[1 + \tfrac{1}{4}\omega_p'^2(\Delta t)^2]^{-1}, \tag{2}$$

in which $\omega_p'^2 = (4\pi e^2/m)[n_e^{(m)} + (m/M)Z^2 n_i^{(m)}]$. For reduced noise with this centering the n_α and j_α terms have been "softened" with the factors $\beta = \gamma = 0.25$, as described in Ref. 10. The particles are then advanced with $du/dt = q_\alpha E^{(*)}/m_\alpha$ and $dx/dt = u$.

First, we examine thermal transport in an end-heated, uniform plasma near the critical conditions for a 1.06-μm laser, Figs. 1(a)–1(d). The density is 10^{21} cm^{-3} and the initial temperature is 0.43 keV ($\equiv T_c$), while $Z = 10$ (for SiO_2). To the right of the vertical fiducial we heat the electrons to 1.72 keV ($\equiv T_h$). The results are for 2.9 ps of heating. We employ 200 particles/cell, and 100 cells with a mean free path $\lambda_{\mathrm{mfp}}(T_c) \equiv a\tau_e = \Delta x$

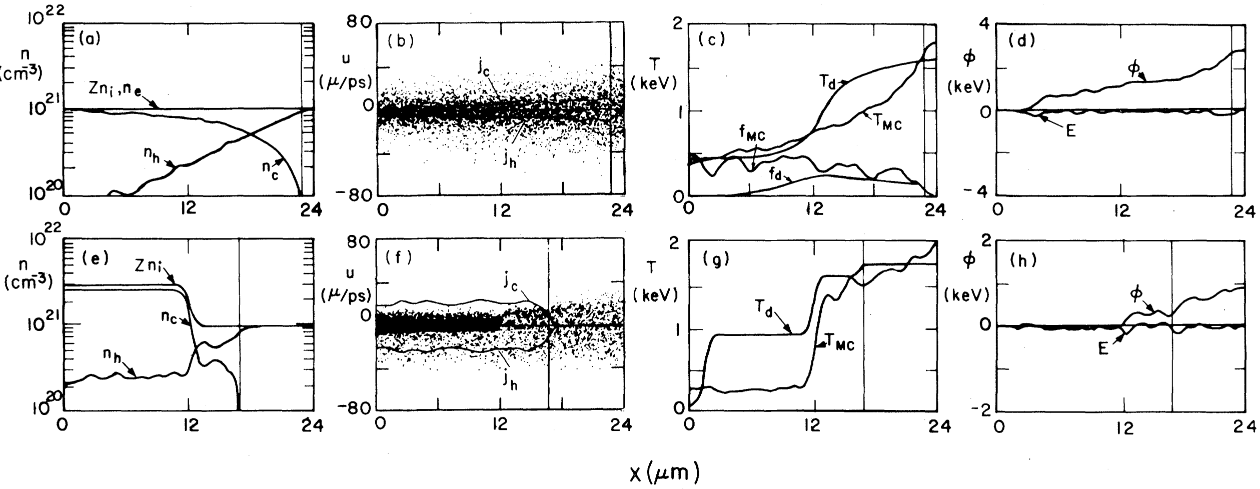

FIG. 1. (a)–(d) An end-heated uniform plasma; (e)–(h) transport through a density gradient. $j_{h,c} \equiv 5\times 10^{-21} n_{n,c} \times u_{h,c}$ (cm^{-3} μm/ps), E in arbitrary units.

$=0.24$ μm; thus, $\lambda_{\mathrm{mfp}}(T_h)=16\Delta x$. Also $\Delta t = 6\times 10^{-3}$ ps. Here $a=(kT/m)^{1/2}$ and τ_e is Braginskii's[11] electron scattering time with $\ln\lambda$ set to 10. For comparison we have conducted a parallel calculation of limited diffusive transport for the same energy deposition, solving $\partial T/\partial t = -\partial q_d/\partial x$, in which $q_d = q_u/[1+|q_u|(fn_e ma^3)^{-1}]$, where $q_u=-K/\gamma(Z)\,\partial T/\partial x$, $K=5n_e T\tau_e/2m$ is Braginskii's conductivity,[11] and $\gamma(Z=10)=0.266$. The Monte Carlo heat flux is $q_{\mathrm{MC}}=0.5mn_e\langle c^2 u\rangle$.

Figure 1(a) shows that $n_e \equiv n_c + n_h \simeq Zn_i$, verifying that the E field in Eq. (2) maintains quasineutrality; cold electrons are given a hot label upon being kicked into the heating region. Frame (d) plots E and $\varphi = -e\int_0^x E\,dx$. Frame (b) shows the electron phase space with superimposed currents $j_h \simeq -j_c$. The electrons are specular at the right boundary and absorbed at the left with current-conserving remission at T_c. Frame (c) plots the temperatures, and the *locally* normalized heat fluxes, $f_{\mathrm{MC},d}=|q_{\mathrm{MC},d}|(mn_e a^3)^{-1}$, for $f=0.3$ in the diffusion calculations.

With this limitation, $T_d \simeq T_{\mathrm{MC}}$ in the heating region, and $f_d \simeq f_{\mathrm{MC}}$ at the heater boundary. With stronger limitation, e.g., for $f=0.1$, the heater T_d climbs to 3.0 keV, while for $f=0.5$ it drops to only 1.3 keV. In these cases the boundary fluxes are mismatched, as well. This fixes f at 0.3. Moreover, when the scatter and drag are "off," similar calculations fix f at 0.5.

For a related problem Bell, Evans, and Nicholas[6] have calculated $f=0.1$. Our value of Debye length per Δx is about equal to Bell's. However, at $\frac{3}{4}T_h$ Bell has $N_{\mathrm{D}}=100$ as the number of particles in his Debye sphere, while for our conditions $N_{\mathrm{D}}=2510$. Thus, Bell's simulation is much more collisional.[12] Indeed, if we bring our λ_{mfp} values into accord with Bell's by lowering our temperatures to $T_h, T_c = 0.32, 0.08$ keV, we find that $f_{\mathrm{MC}} \leqslant 0.1$, and also that $f_d \leqslant 0.1$, even with $f=0.6$, because of the strong collisions. Alternatively, at our $T_{h,c}$ Bell's results could apply[12] to a much larger system.

Since here there is little true inhibition at near-critical conditions, it is significant that $T_d \gg T_{\mathrm{MC}}$ in the *body* of the plasma near the heater. This is due to the nonlinear limited-diffusion front being convex, while the Monte Carlo front is concave, much like a free-streaming expansion. Thus, the diffusion wave heats matter neighboring the heater to relatively higher temperatures, which should lead to relatively higher radiative emissions and increased ablation. The Monte Carlo profile heats more matter in depth to lower temperatures. To bring the two-body temperatures into greater accord, one could increase f, at the expense of a thermal mismatch in the heater region.

This effect is more pronounced and longer lasting in an inhomogeneous plasma. To demonstrate this we raise n_e to 3×10^{21} cm^{-3} for $x<12.1$ μm, while turning off the scatter and drag, and thickening the heating region. Also, we lower T_c to 0.1 keV. Figures 1(e)–1(h) show the results at 3.8 ps. Frames (f) and (g) show that in a two-component system of hot and cold electrons, the hot electrons tend to dominate in the corona, where $T_{\mathrm{MC}} \simeq T_h$. This is essentially coronal *ther-*

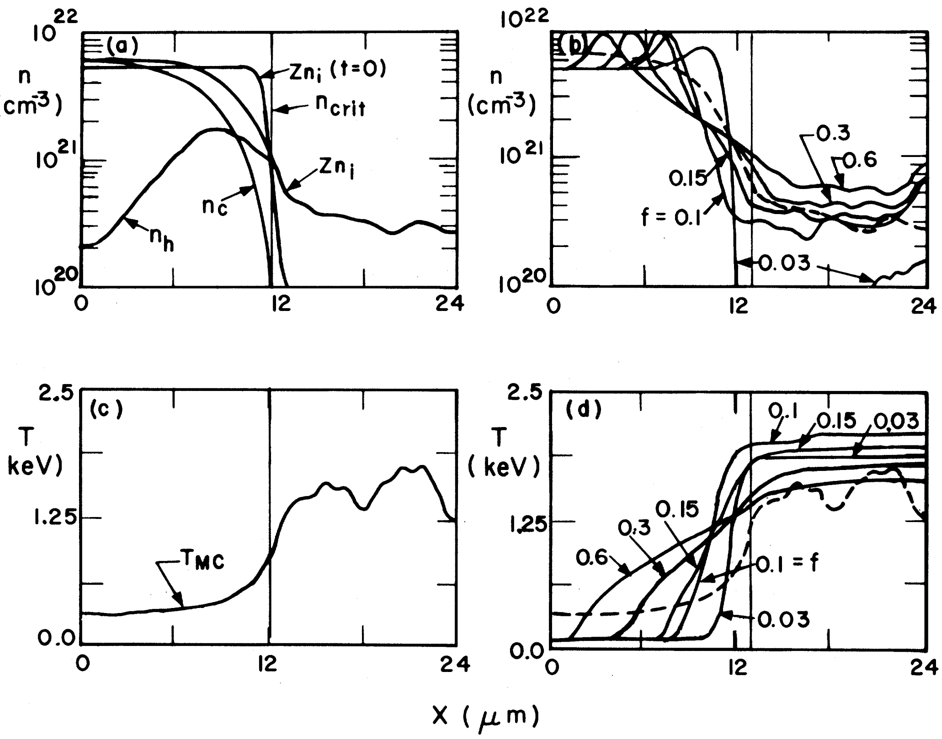

FIG. 2. Ablation calculated (a) and (c) with Monte Carlo, and (b) and (d) via flux-limited single-group diffusion. Solid curve, T_d; dashed curve, T_{MC}. Diffusion critical fiducial is for $f = 0.6$.

mal decoupling.[13,14] In the denser interior there is a mixture of hot and cold electrons, so that for a large density jump and moderate ΔT, $T_{MC} = (n_h T_h + n_c T_c)/n_e \simeq T_c$. Thus, at an interface there can be a significant temperature jump, as in rarefaction shocks,[15] but, with a specular left boundary, for example, no necessity for heat flow. Similarly, the fully flux-limited diffusion profile is pathological, dropping for any f to a plateau temperature, $T_d = T_h(n_h/n_c)^{2/3}$, in order to maintain heat flux continuity. For frames (e)–(h) $f = 0.6$. With progressively larger f the penetration speed of the T_d front beyond the plateau is reduced, holding T_d closer to T_{MC}.

The cold electrons are retained at high density by a potential barrier near $x = 12$ μm of $O(T_c \simeq 0.1$ keV). Frames (f) and (h) show that beyond 12 μm an attractive E field develops to maintain $j_c \equiv n_c u_c$ continuity through the density jump. Here, this results in a $\Delta\varphi = 0.3$ keV potential rise and a net double layer structure.[16] This $\Delta\varphi$ is larger for smaller T_c, and sufficiently moderate Δn such that both u_c, as $x \to 0$, and the cold-tail population are small. Two-stream instability tends to fill in the phase space, and populate the cold tail, reducing $\Delta\varphi$. An additional φ rise, here to $\frac{1}{2}T_h$, is needed to sustain quasineutrality in the heater. These $\Delta\varphi$ produce true inhibition, and with ion motion, a steepening of the expanding density interface, and density depression in the heating region. The restoration of collisions in this problem smears the temperature profiles somewhat, but leaves these conclusions invariant.

Finally, we study the combined consequences of these effects in Fig. 2. The energy is absorbed by inverse bremsstrahlung from a 1.06-μm laser at 10^{15} W/cm^2. Initially, the plasma runs from 5×10^{21} cm^{-3} down to a plateau at 10^{20} cm^{-3}, $T_e = 0.1$ keV, and $Z = 10$. In the limited single-group electron diffusion calculations used for comparison the ions are accelerated by $E = -(en_e)^{-1} \times \partial(n_e T)/\partial x$. The results are for 9.6 ps; with the nucleon/electron mass ratio of 100 employed, this corresponds to 41 ps of physical time. Frame (b) shows that the best match between the Monte Carlo and the diffusion *coronal density* profiles is obtained with $f =$ 0.15 or 0.1. Frame (d) shows a low-temperature Monte Carlo percursor ahead of all the diffusion fronts. However, in the *overdense plasma near critical*, $T_d \to T_{MC}$ as $f \to 0.03$ or lower. Alternatively, the *coronal temperatures* are best matched with $f = 0.6$. This is consistent with the apparent need for separate internal and coronal flux limiters in the recent transport and absorption experiments at Ecole Polytechnique.[17] Thus, the need for a severe transport flux limiter has been traced to deficiencies in the diffusion modeling. Experimental temperature profile determinations near critical should help to verify this conclusion.

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