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Routes to irreversibility in collective laser–matter interaction

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Abstract

Collective absorption, so far determined by numerical simulations, is explained in physical terms for cold and warm plasma. After deducing a few general relations for flat targets, interface phase mixing and nonadiabatic electron acceleration in the skin layer are identified as the main physical processes leading to irreversibility, that is, to absorption.

1. THE PROBLEM

Collisional absorption very soon becomes inefficient when a superintense laser pulse impinges on condensed matter owing to the inverse third power dependence of the electron-ion collision frequency on the quiver velocity, $\nu_{ei} \sim v_{\text{os}}^{-3}$. Further light conversion has to rely on collective, that is, collisionless absorption. Hence, collective absorption is a key issue in superintense laser–solid interaction. At present there essentially exist four models, independent of each other, which in chronological order are the so-called $\boldsymbol{j} \times \boldsymbol{B}$ heating (Kruer & Estabrook, 1985), Brunel effect (Brunel, 1987, 1988), nonlinear skin layer absorption (Mulser *et al.*, 1989), and vacuum heating (Gibbon & Bell, 1992). They are distinct from each other and each of them is commonly viewed as contributing to collective absorption of superintense laser beams.

Depending on target density and interaction geometry (e.g., oblique incidence), and on personal preference, often a single one out of the four models is favored. No general analysis exists up to now of their real relevance in specific interaction experiments. On the other hand, understanding the underlying physics of the collective absorption process may be vital for optimizing a desired interaction process, for example, energetic electron production and efficient hard X-ray generation.

The idea of $\boldsymbol{j} \times \boldsymbol{B}$ heating, as originally proposed, is the following. A laser beam impinging perpendicularly onto a target in positive x direction drives a current density $\boldsymbol{j}$ of periodicity 2ω by the Lorentz force $-e\boldsymbol{v}_{\text{os}} \times \boldsymbol{B} = -e(\boldsymbol{E}_0 + \boldsymbol{E}_d)$. The secular field $\boldsymbol{E}_0$ gives rise to the ponderomotive force, whereas the time-dependent component $\boldsymbol{E}_d$ drives $\boldsymbol{j}$ in the direction of the incident beam and does work on it per unit time of magnitude $\boldsymbol{j}\boldsymbol{E}_d$, which when cycle-averaged in a fixed volume element, was supposed to yield a positive contribution, $\overline{\boldsymbol{j}\boldsymbol{E}_d} > 0$. Subsequently, $\boldsymbol{j} \times \boldsymbol{B}$ heating was frequently addressed to explain collective absorption in particle-in-cell (PIC) simulations (Denavit, 1992; Wilks *et al.*, 1992; Wilks, 1993). Currently $\boldsymbol{j} \times \boldsymbol{B}$ heating is well established as a valuable absorption mechanism among theoreticians and experimentalists. However, so far no attempt has been made to clarify by which physical effect *irreversibility*, needed to produce absorption, comes into play, neither in $\boldsymbol{j} \times \boldsymbol{B}$ nor in vacuum heating. The occurrence of an irreversible element in the $\boldsymbol{j} \times \boldsymbol{B}$ heating model was rather taken for granted by the appearance of about 10% absorption in 1-D PIC simulations at normal incidence (Kruer & Estabrook, 1985). It was confirmed, meanwhile, by Vlasov and PIC simulations, that the Brunel mechanism contributes by a small faction only to collective absorption. Finally, we shall see that skin layer absorption is significant. The basic idea of how irreversibility in this mechanism comes into play was outlined correctly for the first time by Mulser *et al.* (1989). However, in detail the mechanism is more subtle, owing to the conservation of canonical momentum in the laser field direction at normal incidence.

To make the basic physics clear, we treat the problem in the most symmetric, that is, 1-D, configuration. First we derive a few very useful general relations. Then we show by which kind of symmetry breaking irreversibility is introduced in a cold plasma and give an estimate of the upper limit of Brunel absorption. In Section 4 we show the route into irreversibility originating from the skin layer in the hot plasma.

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2. GENERAL RELATIONS IN 1-D CONFIGURATION

We assume a plane wave $E_y(x,t)$ normally incident onto a plane plasma surface. The ions are held fixed with a discontinuous transition from vacuum to a highly supercritical (e.g., solid-state) density $n_0 > n_c$. Oblique incidence of a plane wave $\boldsymbol{E}$ under the angle α is reduced to perpendicular incidence by applying a Lorentz boost $v_0 = c\sin\alpha$ in the y direction of the wave vector component $k_y = k\sin\alpha$. With the Lorentz factor $\gamma = (1 - v_0^2/c^2)^{-1/2} = (\cos\alpha)^{-1}$, the following relations for the field quantities in the moving frame (primes) are easily confirmed for p-polarization:

$$E'_y = E_y = \frac{E}{\gamma}, \quad B'_z = \frac{E}{c\gamma}, \quad B'_x = B'_y = 0; \qquad E = |\boldsymbol{E}|, \tag{1}$$

and for s-polarization:

$$E'_y = 0, \quad E'_z = \frac{E_z}{\gamma}, \quad E_z = E, \quad B'_y = B_y = \frac{B}{\gamma},$$

$$B'_x = 0; \qquad B = |\boldsymbol{B}| = \left|\frac{\boldsymbol{k}}{\omega} \times \boldsymbol{E}\right|, \tag{2}$$

by making use of the self-evident transformations

$$E'_\parallel = E_\parallel, \quad E'_\perp = \gamma(\boldsymbol{E} + \boldsymbol{v}_0 \times \boldsymbol{B})_\perp, \quad B'_\parallel = B_\parallel,$$

$$B'_\perp = \gamma\left(\boldsymbol{B} - \frac{\boldsymbol{v}_0}{c^2} \times \boldsymbol{E}\right)_\perp.$$

The transformation law of $B'_\perp$, for instance, follows from that of $E'_\perp$ by performing the identity boost first with $\boldsymbol{v}_0$ and then with $-\boldsymbol{v}_0$. Of course, Eqs. (2) immediately follow in the four-vector notation from the electromagnetic field tensors.

In addition

$$k'_x = k' = k_x = \frac{k}{\gamma}, \quad \omega' = \frac{\omega}{\gamma}, \quad n'_0 = \gamma n_0, \quad j_0 = \gamma e(n_0 - n_e)v_0 \tag{3}$$

follows. No ambiguity arises if the prime sign ($'$) is omitted from here on for simplicity. To treat p- and s-polarization in a unified way, in s-polarization the boosted frame is rotated around the x' axis by $+\pi/2$ and $v_0 = c\sin\alpha = -\dot{z}$ is used. Introducing the vector and scalar potentials $A_y(x,t), \phi(x,t)$ in the boosted frame, with

$$E_y = -\frac{\partial}{\partial t}A_y, \quad B_z = \frac{\partial}{\partial x}A_y, \quad E_x = -\frac{\partial\phi}{\partial x},$$

the relevant Maxwell and Lorentz equations read

$$-\varepsilon_0 c^2 \frac{\partial}{\partial x}B_z = \varepsilon_0 \frac{\partial}{\partial t}E_y + j_y \tag{4}$$

$$0 = \varepsilon_0 \frac{\partial}{\partial t}E_x + j_x \tag{5}$$

$$\frac{\partial E_x}{\partial x} = \frac{e}{\varepsilon_0}(n_0 - n_e) \tag{6}$$

$$\frac{d}{dt}mv_y = e\frac{d}{dt}A_y \tag{7}$$

$$\frac{d}{dt}mv_x = e\frac{\partial\phi}{\partial x} - ev_y\frac{\partial A_y}{\partial x}. \tag{8}$$

All equations are relativistically correct; $m = \gamma m_e$. For simplicity, the ion charge number is assumed to be unity.

Now assume that (1) a steady state is reached and that (2) E_x is periodic, with periodicity $n\omega$, $n \geq 1$, or combinations of it. Poynting's theorem states

$$\frac{\partial}{\partial x}S + j_x E_x + j_y E_y = 0.$$

From Eq. (5) follows

$$j_x E_x = -\frac{\varepsilon_0}{2}\frac{\partial}{\partial t}E_x^2. \tag{9}$$

Cycle-averaging leads to the conclusion

$$\overline{j_x E_x} = 0. \tag{10}$$

Under the above assumptions, it is of general validity regardless of whether the electron temperature T_e is zero or finite. In particular, it holds in steady-state resonance absorption in flat density gradients. The relation may serve as a test for the accuracy of a numerical simulation. In wakefield acceleration it was shown explicitly by PIC simulations that $\overline{j_y E_y}$ is much larger than $\overline{j_x E_x}$ (Gahn *et al.*, 1999). This is not surprising since the configuration is close to 1-D and a quasistationary state builds up. Assumption (1) is well fulfilled in 1-D Vlasov simulations and Eq. (10) holds (e.g., Ruhl & Mulser, 1995, Figure 9). The field E_x is the sum of the driver field $E_d = v_y B_z$ and the induced field E_{in} obeying $\partial E_{in}/\partial x = \partial E_x/\partial x$ of Eq. (6), $E_x = E_{in} + E_d$. Hence, Eq. (10) implies

$$\overline{j_x E_d} = -\overline{j_x E_{in}}.$$

With zero electron temperature and immobile ions, energy convection is absent in ideal fluid dynamics. Thus, $\overline{j_x E_{in}}$ vanishes. On the other hand $\overline{j_x E_d}$ is the cycle-averaged irreversible work done by the $\boldsymbol{j} \times \boldsymbol{B}$-force density. Therefore, in the absence of discontinuities, one concludes from fluid

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dynamics that under conditions (1) and (2) no $\boldsymbol{j} \times \boldsymbol{B}$ heating occurs at $T_e = 0$. This conclusion is still valid when, in a steep continuous transition from vacuum to high density, a plasma oscillation is resonantly excited. At finite electron temperature, $\overline{j_x E_x} = 0$ still holds under (1) and (2); however, $\overline{j_x E_d} = 0$ can no longer be inferred since now, for example, the high harmonics of the electron plasma wave carry energy into the dense target. In case of $n_0 \gg n_c$, the energy transported away is expected to be very small. The conclusion is that if collisionless absorption is high, which is the case at large angle of incidence α (Ruhl & Mulser, 1995), that is, v_0 is large, it cannot originate from $j_x E_x$ in the boosted frame. High collisionless absorption (up to 70%) is proved experimentally (Feurer *et al.*, 1997).

The only alternative is absorption by $\overline{j_y E_y}$. Multiplying Eq. (4) by $E_y = -\partial_t A_y$ and substituting B_z by $\partial_x A_y$ yields

$$j_y E_y = \varepsilon_0 c^2 \frac{\partial A_y}{\partial t} \frac{\partial^2 A_y}{\partial x^2} - \frac{\varepsilon_0}{2} \frac{\partial}{\partial t} \left(\frac{\partial A_y}{\partial t} \right)^2.$$

In the cycle-averaging process, the second term on the right-hand side vanishes under conditions (1) and (2). Hence,

$$\overline{j_y E_y} = \varepsilon_0 c^2 \overline{\frac{\partial A_y}{\partial t} \frac{\partial^2 A_y}{\partial x^2}}. \tag{11}$$

A third important relation results from Eq. (7),

$$m v_y = e A_y + \theta m_0 v_0, \quad m_0 = \gamma_0 m_e, \quad \gamma_0 = \cos^{-1} \alpha. \tag{12}$$

For p-polarization $\theta = 1$; for s-polarization $\theta = 0$. It expresses the conservation of the canonical momentum $p_y = m v_y - e A_y$ owing to translational invariance along y in the boosted frame. In other words, in a stationary 1-D plasma, thermal electrons ($v_x \neq 0$) cannot carry transverse oscillation energy out of the skin layer regardless of how thin it is.

A salient signature of intense laser–matter interaction is the generation of energetic electrons. To establish a link between this phenomenon and absorption, Eq. (11), the momentum balance may be used (here nonrelativistic for simplicity),

$$\frac{\partial}{\partial t} \frac{\rho_e}{2} v_i^2 + \frac{\partial}{\partial x_j} \rho_e v_i v_j + \frac{\partial}{\partial x_j} p_{ij} = -n_e e [E_i + (\boldsymbol{v} \times \boldsymbol{B})_i];$$

$$p_{ij} = \rho_e \langle (u_i - v_i)(u_j - v_j) \rangle.$$

Hence, under conditions (1)and (2)

$$\overline{j_x E_x} = \overline{v_x \frac{\partial}{\partial x} \rho_e v_x^2} + \overline{v_x \frac{\partial p_{xx}}{\partial x}} + e n_0 v_0 \bar{E}_y = 0, \quad \rho_e = m_e n_e;$$

$$\overline{j_y E_y} = \frac{1}{2} \frac{\partial}{\partial x} \overline{\rho_e v_x \left(\frac{e}{m_e} A_y + \theta v_0 \right)^2} + \overline{\left(\frac{e}{m_e} A_x + \theta v_0 \right) \frac{\partial p_{xy}}{\partial x}}. \tag{13}$$

The total absorption is obtained by spatial integration over the interaction region,

$$\int \overline{j_y E_y} \, dx = \int \overline{\left(\frac{e}{m_e} A_y + \theta v_0 \right) \frac{\partial p_{xy}}{\partial x}} dx. \tag{14}$$

At first glance, the somewhat surprising consequence is that at $T_e = 0$, the term $j_y E_y$ does not lead to absorption. Together with Eq. (10), this seems to contradict all PIC simulations performed until recently and, to a minor extent, corresponding Vlasov modeling of collective absorption. However, one must bear in mind that whenever there is a spread in the electron energy distribution, the pressure tensor p_{xy} and the kinetic temperature T_e do not vanish.

3. INTERFACE PHASE MIXING

At the beginning of laser incidence, $T_e = 0$ is assumed for clarity. Such a simplification is made also in standard PIC simulations, although, in the real experiment, there is some kinetic temperature already introduced by the field ionization process itself (Ruhl, 1994; Mulser *et al.*, 1998) and by inverse bremsstrahlung absorption. According to the foregoing section, an effect is needed for absorption, which adds an irreversible element to the ideal collective cold fluid dynamics. At $T_e = 0$ there exist only two such elements: cold wavebreaking (Mulser *et al.*, 1982; Koch & Albritton, 1974) and symmetry breaking of the $\boldsymbol{j} \times \boldsymbol{B}$ force-driven oscillatory motion. Wavebreaking in a cold fluid is well understood as overlapping and subsequent interpenetration of different fluid elements. A simple estimate below will reveal that this phenomenon occurs only at very high intensities and does generally not come into play in our context. Irreversibility rather starts from the vacuum plasma interface where a fluid element, when crossing this boundary, is subject to different kinds of motion. To make this clear, normal incidence of the laser beam onto the plasma surface at $x = 0$ in the boosted frame is considered. Under the influence of the driven E_d and the induced field E_{in}, a fluid element at its original position $x_0 > 0$ undergoes a displacement $\delta(x_0, t)$. Multiplying Eq. (6) by the electron velocity $v_x = \dot{\delta}$ and adding to Eq. (5) yields

$$\frac{d}{dt} (\varepsilon_0 E_x - \theta e n_0 \delta) = \frac{d}{dt} (\varepsilon_0 E_x + \theta P) = 0, \quad P = -e n_0 \delta. \tag{15}$$

Inside the plasma, that is, $x_0 + \delta \geq 0$, is $\theta = 1$; for $x_0 + \delta < 0$ holds $\theta = 0$. Combining this with the total time derivative of Eq. (8) in nonrelativistic form leads to

$$\frac{d^3}{dt^3} \delta + \theta \omega_{p0}^2 \frac{d}{dt} \delta = -\frac{e}{m_e} \frac{d}{dt} E_d. \tag{16}$$

The equation is strictly valid as long as no electron fluid elements interchange, that is, when from $x_{01} < x_{02}$ follows $x_1 < x_2$ for the actual positions $x = x_0 + \delta(x_0, t)$. If Eq. (16)

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is integrated once, for $\theta = 1$ the displacement δ is governed by

$$\ddot{\delta} + \omega_{p0}^2 \delta = -\frac{e}{m_e} E_d; \qquad x_0 + \delta \geq 0. \tag{17}$$

For $\theta = 0$, the integration yields

$$\ddot{\delta} - G(x_0) = -\frac{e}{m_e} E_d, \quad G = \omega_{p0}^2 x_0; \qquad x_0 + \delta < 0.$$

In analogy in Eq. (17), this can be cast into the form

$$\ddot{\delta} + \omega_p^2 \delta = -\frac{e}{m_e} E_d, \quad \omega_p^2 = -\omega_{p0}^2 \frac{x_0}{\delta}. \tag{18}$$

The beauty of Eq. (17) is that, although it is highly nonlinear in (x, t), it becomes linear with constant plasma frequency in the Lagrangean variables (x_0, t). Equation (18) contains instead the free fall term G. In oscillating across the boundary, an electron fluid element switches from the constant eigenfrequency ω_{p0} to the continuously varying one ω_p. Therefore, when oscillating back to the interface, it has accumulated a phase difference relative to the oscillation mode in the interior which results in a drift in addition to its regular oscillation, in perfect analogy to a collision. This is the underlying physics of irreversibility leading to collisionless absorption in 1-D at $T_e = 0$. It is nonstationary, therefore violating condition (1). With $E_d = \hat{E}_d \sin \Omega t$ switched on adiabatically the solutions of (17) and (18) are

$$x_0 + \delta \geq 0: \quad \delta(x_0, t) = -\hat{\delta} \sin \Omega t; \qquad \hat{\delta} = \frac{e}{m_e(\omega_{p0}^2 - \Omega^2)} \hat{E}_d$$

$$x_0 + \delta < 0: \quad \delta = \delta_V(x_0, t) = \hat{\delta}_V (\sin \Omega t - \sin \Omega t_0)$$

$$+ \frac{G}{2}(t - t_0)^2 - \Omega(t - t_0)(\hat{\delta}_V + \hat{\delta}) \cos \Omega t_0;$$

$$\dot{\delta}_V(x_0, t) = \Omega \hat{\delta}_V \cos \Omega t + G(t - t_0) - \Omega(\hat{\delta}_V + \hat{\delta}) \cos \Omega t_0;$$

$$\hat{\delta} \sin \Omega t_0 = x_0; \qquad \hat{\delta}_V = \frac{e}{m_e \Omega^2} E_d. \tag{19}$$

Thereby the approximation $\hat{E}_d(x_0 + \delta) \simeq \hat{E}_d(x_0)$ was made, that is, harmonics are suppressed. Depending on the starting position x_0, there are particles which, when falling back from the vacuum, cross the boundary at high velocity $\dot{\delta}_V$. In the interior they proceed at nearly constant speed because their motion is neutralized by a slow back current of the bulk plasma. To give an example, an electron entering the vacuum at $\Omega t_0 = 80°$ crosses the boundary at $\Omega t = 217.1°$ with speed $\dot{\delta}_V / \Omega \hat{\delta}_V = 1.52$; it corresponds to a kinetic energy 2.3 times its oscillation energy in vacuum.

To show in detail how interface phase mixing works and how irreversibility originates, a cold 1-D PIC code has been written just for this purpose which is free from self-heating effects of standards PIC codes (Steinmetz, 2000). In Figure 1, a laser beam of irradiance $I\lambda^2 = 10^{19}$ Wcm^{-2} μm^2 is incident from the left-hand side onto a semi-infinite target of density $n/n_c = 10$ under $\alpha = 0°$, 45°, and 60°. The interface is placed at $x = 0$. The plots show the time history of the first

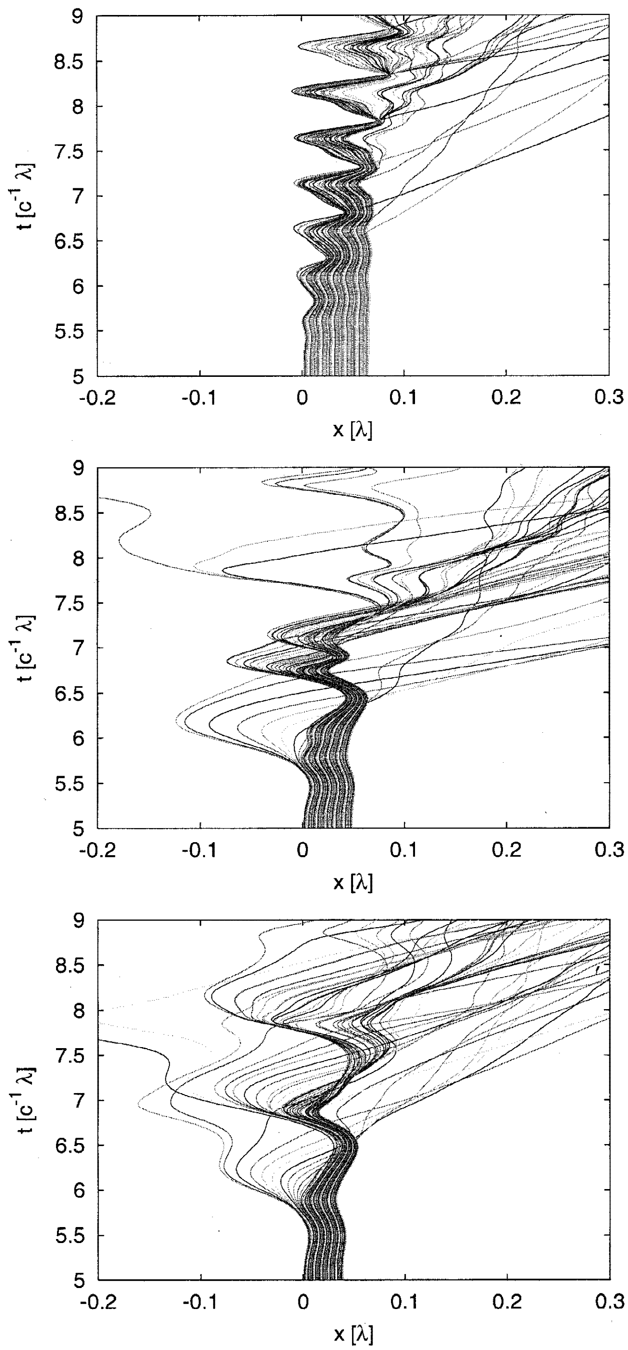

Fig. 1. Interface phase mixing. Particle-in-cell (PIC) simulation of cold plasma heating ($n_0 = 10n_c$) by a laser of irradiance $I\lambda^2 = 10^{19}$ Wcm^{-2} μm^2. The laser beam is turned on over three cycles; its front reaches the target surface after five cycles. The angle of incidence from top to bottom is $\alpha = 0°, 30°, 60°$. The pictures show the dynamics of 20 layers (positive x in units of laser wavelength λ) as a function of time t (in units of laser cycles).

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50 layers under the action of the driving field E_d and the induced self-consistent field E_{in}. Under normal incidence (top) the driver oscillates at $\Omega = 2\omega$ owing to $E_d = v_y B_z$; in the second and third pictures, E_d is a combination of components $v_0 B_z \sim \omega$ and $v_{os} B_z \sim 2\omega$. It is largest for $\alpha = 60°$. It is impressive to observe how strong the interface phase mixing develops in the last picture: a small phase delay relative to the regular motion of the skin layer leads to crossing and subsequently to an unbound motion into the target. The vacant position at the surface is filled by a slow return current, establishing in this way the required quasi-neutrality. In the vacuum, the layers follow Eq. (19) approximately, because the few layer crossings there do not substantially alter the backdriving force $G(x_0) = \omega_{p0}^2 x_0$ in Eq. (18). It is easily seen that crossing events introduce an additional irreversible element.

Some conclusions extracted from the numerical simulation may shed additional light on collective absorption:

1. Only the region $n_e < n_c$ contributes to cycle-averaged absorption $\overline{j_y E_y}$ (Fig. 2); $\overline{j_x E_x} = 0$ is well satisfied in the boosted frame, as expected;
2. From Figure 2, Brunel absorption $(x < 0)$ amounts to 25% of total absorption;
3. If in the simulation the layers are reflected when they reach the vacuum boundary $x = x_0 + \delta(x_1, t_0) = 0$, instead of moving out into the vacuum, the absorption increases slightly.

These facts are of particular relevance. They tell us that interface phase mixing is an important phenomenon and that the Brunel mechanism does not significantly contribute to total absorption. On the other hand, increased absorption under geometrical reflection is understandable because, by reflection, dephasing is increased for those layers not moving far out into vacuum because, at a reflecting boundary, a layer has no chance to undergo a quasiadiabatic phase shift.

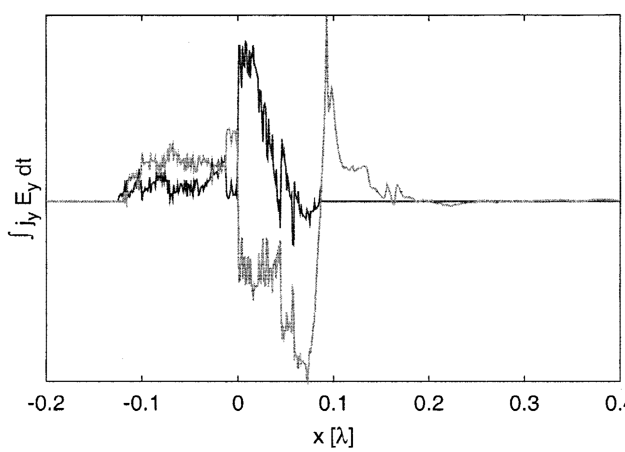

Fig. 2. Cycle averaged absorption $\int j_y E_y dt = \overline{j_y E_y}$ of a $I\lambda^2 = 10^{19}$ Wcm^{-2} μm^2 laser beam under $\alpha = 30\%$ of incidence onto $n_0 = 10 n_c$. Bold curve: $\overline{j_y E_y}$ of the underdense electron region (rarefaction is due to dynamics); light curve: $\overline{j_y E_y}$ in $n_e > n_c$. Overall absorption is 23–24% and is entirely due to underdense electron region. The fraction of Brunel absorption $(x < 0)$ amounts to $\frac{1}{4}$ of total absorption.

For cold wave breaking to occur, $\partial\delta/\partial x_0 = -1$, that is, $|\partial\hat{\delta}/\partial x_0| = 1$, must hold (Mulser *et al.*, 1982). At normal incidence, $E_d = v_y B_z$ is determined from Eq. (18) and $B_z = \partial A_y/\partial_x$,

$$E_d = \frac{1}{2}\frac{e}{m_e}\frac{\partial A_y^2}{\partial x}.$$

In the strongly overdense skin layer $\hat{A}_y(x) = \hat{A}_0 e^{-\kappa x}$, $\kappa = k\omega_{p0}/\omega$ can be set and hence equation (19) yields

$$\left|\frac{\partial\hat{\delta}}{\partial x_0}\right| = 2\kappa^2 \frac{e^2}{m_e^2(\omega_{p0}^2 - 4\omega^2)} A_y^2 \simeq 2\frac{e^2 A_y^2}{m_e^2 c^2}, \quad \omega^2 \ll \omega_{p0}^2. \quad (20)$$

This becomes unity at about $I\lambda^2 = 10^{18}$ Wcm^{-2} μm^2 only. In Figure 1, $I\lambda^2$ has exceeded this limit already at six laser periods and no indication of cold wavebreaking can be observed. In another simulation (Fig. 3), 200 particle orbits are shown with the same laser pulse under $\alpha = 30°$ of incidence with the neat result that wavebreaking is absent; all absorption is due to phase mixing and, to a much weaker degree, Brunel acceleration.

4. NONLINEAR SKIN LAYER ABSORPTION

When the plasma electrons have heated up to a kinetic temperature of several kiloelectronvolts, a quasiequilibrium is established between its pressure p_{ij}, radiation pressure, and induced thermoelectric field on the secular, that is, slow, time scale. In this environment, the electrons can be assumed to move with their random ("thermal") velocities v_e across the skin layer. At the plasma–vacuum interface, we assume that they are geometrically reflected. This is a good approximation because the Debye layer of thickness λ_D has a much more reduced extension than the skin layer $\delta_s = 1/\kappa$.

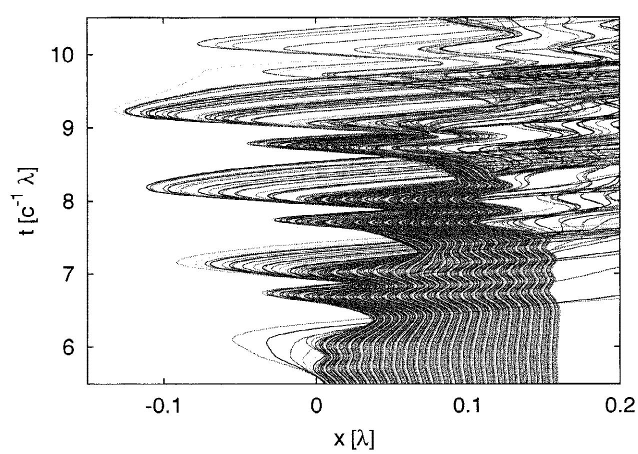

Fig. 3. Dynamics of 200 cold plasma layers induced by the laser pulse of Figure 1 under $\alpha = 30°$ of incidence. No indication of cold wavebreaking.

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An electron starting at $x_0 > \delta_s$ with a velocity v_{0x} towards the vacuum boundary is subject to the Lorentz force $f_L = -e(v_{\mathrm{os}} - v_0)B_z$ with B_z having the form

$$B_z(x,t) = \frac{\kappa}{\omega}\hat{E}e^{-\kappa x}\sin\omega t, \quad x = x_0 + v_{0x}t + \int^t\int^{t'} f_L\,dt''dt'. \tag{21}$$

An electron moving at low speed may undergo several oscillations during crossing the skin layer. In this case the particle has almost no energy gained after arriving again at x_0. However, at $v_{0x} \simeq c$, the phase of B_z changes only slightly, that is, $\Delta\varphi \simeq (2\delta_s/c)\omega = 4\pi\omega/\omega_p \ll 2\pi$, and the electron is subject to a nearly static accelerating and decelerating force. There exists an optimum velocity v_{0x} for gaining energy. To see this in detail one may move in such a parameter region ("temperature" T_e, irradiance $I\lambda^2$, and angle of incidence α) that the position is approximately given by $x = x_0 + v_{0x}t$. If then $f_L \simeq -ev_0 B_z$ is negative before reflection and positive after reflection, the gain in velocity v_{0x} and energy $\mathcal{E}_e$ is

$$\Delta v_{0x} \simeq ev_0\frac{B_z}{2}\frac{\delta_s}{v_{0x}} > 0, \qquad \Delta\mathcal{E}_e = m_e[v_{0x}\Delta v_{0x} + \tfrac{1}{2}(\Delta v_{0x})^2].$$

This process may happen with equal likelihood at the opposite phase, with an energy loss $\Delta\mathcal{E}_e = m_e[-v_{0x}\Delta v_{0x} + (\Delta v_{0x})^2/2]$, and hence in this special case the net energy gain is $m_e(\Delta v_{0x})^2$. The general case is characterized by no momentum gain over the ensemble average, $\langle\Delta m_e v_{0x}\rangle = 0$; however in the single event $\Delta m_e v_{0x} \neq 0$ ($\Delta m_e v_{0x} = 0$ occurs very seldom). This implies the positive energy gain

$$\langle\Delta\mathcal{E}_e\rangle = \tfrac{1}{2}m_e[\langle(v_{0x} + \Delta v_{0x})^2\rangle - \langle v_{0x}^2\rangle] = \tfrac{1}{2}m_e\langle(\Delta v_{0x})^2\rangle \tag{22}$$

after averaging over all phases. The numerical evaluation of the exact integrals shows that the absorption by this mechanism can become as high as 40% of the total power incident under large angles α on flat targets (Ruhl, 1994). In light of Section 2, Eqs. (10) and (14), we conclude on the basis of Eq. (22) that the $\boldsymbol{v}_y \times \boldsymbol{B}_z$ force in the skin layer introduces an irreversible element leading to an absorptive dephasing between j_y and E_y, which in terms of a fluid description has its origin in the pressure tensor p_{ij}.

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