Radiation Damping Effects on the Interaction of Ultraintense Laser Pulses with an Overdense Plasma

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(Received 24 May 2001; published 19 April 2002)

A strong effect of radiation damping on the interaction of an ultraintense laser pulse with an overdense plasma slab is found and studied via a relativistic particle-in-cell simulation including ionization. Hot electrons generated by the irradiation of a laser pulse with a radiance of $I\lambda^2 > 10^{22}$ W μm^2/cm^2 and duration of 20 fs can convert more than 35% of the laser energy to radiation. This incoherent x-ray emission lasts for only the pulse duration and can be intense. The radiation efficiency is shown to increase nonlinearly with laser intensity. Similar to cyclotron radiation, the radiation damping may restrain the maximal energy of relativistic electrons in ultraintense-laser-produced plasmas.

DOI: 10.1103/PhysRevLett.88.185002 PACS numbers: 52.38.–r, 52.65.Rr

The recent advance of the high-power, short pulse laser technique allowing laser intensities over 10^{20} W/cm^2 [1] makes the study of relativistic effects in the interaction of such laser pulses with high-Z matter interesting for various applications [1–3]. With a laser pulse energy of about 1 kJ [1] and a pulse duration of 20 fs, the laser intensity can exceed 10^{22} W/cm^2 which is critical for well known relativistic effects such as radiation damping [4–10].

The electron self-interaction problem, radiation damping (RD), is central to both classical and quantum electrodynamics [5,6]. Usually, this effect is attributed to the physics of relativistic electron beams [4,5]. In ordinary plasmas, the effect of RD is negligible because the plasma field usually is not strong enough for radiation to dominate the particle motion. However, in overdense plasmas, the RD can become important due to Thomson scattering at laser intensities over 10^{22} W/cm^2 which corresponds to the normalized laser amplitude of $a_0 = eE_0/mc\omega \sim 50$ where E_0 is the amplitude of the electromagnetic wave and ω is the laser frequency. Roughly, this can be estimated through the comparison of the RD force and the Lorentz force. The maximal energy of an electron upon the laser radiation, $\gamma \sim a_0^2/2$, with $a_0 = eE_0/mc\omega$, can be found by using the solution for motion of an electron with zero initial velocity in the plane wave as given in Ref. [6]. To prevent the expulsion of electrons from the skin layer, a potential difference $\Delta\Phi$ should appear with $e\Delta\Phi \sim \gamma$. The radiation damping force that acts upon electrons accelerated by the potential difference, $f_{\mathrm{RD}} \sim 2e^2\omega^2\gamma^2 a_0^2/3c^2$ [9], could be comparable with the Lorentz force, $f_{\mathrm{L}} = mc\omega a_0$, at $a_0 = (6mc^3/e^2\omega)^{1/5} = [(3/\pi)\lambda/r_e]^{1/5}$ where r_e is the classical electron radius and λ is the laser wavelength. The expression within the parentheses is a very large number; however, with the fifth root this gives $a_0 = 50$. In the simple plane wave, the RD effects are very small because a free electron accelerated by the wave does not radiate. In contrast to this, in overdense plasmas of multiple-charged ions with $\omega_{\mathrm{pl}}^2/\omega^2\gamma > 1$, which are opaque for the laser light, a strong longitudinal electric field is produced to compensate the laser pressure. This field can accelerate plasma electrons up to ultrarelativistic energies in the direction counter to both the incident and reflected parts of the laser light with consequent Thomson scattering [11]. That may instigate strong radiation damping converting the laser energy to plasma radiation and affecting the electron energy distribution and ion acceleration. Moreover, the laser light coming into the plasma is not a simple plane wave, and the RD force can be very strong for a laser field over $a_0 > 50$. In this Letter, we consider the effect of RD on the interaction of ultraintense laser pulses with overdense plasmas. Although finally the radiation is a result of Thomson scattering by a relativistic electron, the whole effect is due to collective plasma motion and not single particle motion in the laser field. This radiation can be of short duration, short wavelength, and high intensity.

A general solution of the Maxwell and Lorentz equations should contain both external fields and fields radiated by particles and their effect on the particle motion. However, the calculation of the full effect of the radiation requires a three-dimensional simulation for spherical waves emitted by a particle with high spatial resolution, $\Delta x \ll \lambda/2\pi$, where λ is the wavelength of the radiation. In the very strong relativistic case, the characteristic wavelength of the radiation can be estimated as $\lambda \sim \lambda_0/\gamma^2 a_0$ ($\lambda/2\pi \sim 10^{-10}$ cm for $a_0 = 50$ and $\lambda_0 = 1$ μm). This makes any direct calculation of the radiation impractical. As a result, we use a one-dimensional relativistic particle-in-cell simulation including radiation damping in the frame of the Dirac-Lorentz equation [9,12] with the RD force that is exact for a point particle [13]. Since we use simulation cells which are greater than λ, this approximation is also limited by the condition $N_e\lambda^3 \ll 1$, where N_e is the electron density, which assumes the absence of coherent radiation which could be proportional to N_e^2. For $N_e \sim 10^{24}$ cm^{-3} approximation is correct at λ less than 1 Å which is valid under the condition of strong RD.

In a one-dimensional particle-in-cell calculation with numerical cells of size greater than the radiation wavelength, we use the Dirac-Lorentz equation [8,9,12]

$$mc\,\frac{du^i}{ds} = \frac{e}{c}F^{ik}u_k + g^i, \qquad (1)$$

where F^{ik} is the well known electromagnetic tensor [8], $\boldsymbol{u}$ is the velocity four-vector, and g^i is the radiation damping force

$$g^i = \frac{2e^2}{3c}\left(\frac{d^2u^i}{ds^2} - u^i u^k \frac{d^2u_k}{ds^2}\right). \qquad (2)$$

To avoid self-acceleration, g^i is expressed through the electromagnetic fields replacing the velocity $\boldsymbol{u}$ by a solution of Eq. (1) with $g = 0$ [9]. Finally, the radiation damping force acquires the following form, which is exact for a point particle [13]:

$$\boldsymbol{f} = \frac{2e^3}{3mc^3}\gamma^{1/2}\left\{\left(\frac{\partial}{\partial t} + (\boldsymbol{\nu}\cdot\nabla)\right)\boldsymbol{E} + \frac{1}{c}\left[\boldsymbol{\nu}\times\left(\frac{\partial}{\partial t} + (\boldsymbol{\nu}\cdot\nabla)\right)\boldsymbol{H}\right]\right\} + \frac{2e^4}{3m^2c^4}\left\{\boldsymbol{E}\times\boldsymbol{H} + \frac{1}{c}\boldsymbol{H}\times(\boldsymbol{\nu}\times\boldsymbol{H}) + \frac{1}{c}\boldsymbol{E}(\boldsymbol{\nu}\cdot\boldsymbol{E})\right\} - \frac{2e^4}{3m^2c^5}\gamma^2\boldsymbol{\nu}\left\{\left(\boldsymbol{E} + \frac{1}{c}\boldsymbol{\nu}\times\boldsymbol{H}\right)^2 - \frac{1}{c^2}(\boldsymbol{\nu}\cdot\boldsymbol{E})^2\right\}, \qquad (3)$$

where $\boldsymbol{\nu}$ is the electron velocity, and $\boldsymbol{E}$ and $\boldsymbol{H}$ are the self-consistent electric and magnetic fields in the plasma. The third term is the largest for a relativistic electron. However, at $\nu = c$, for a plane wave $E_y = H_z$, $E_z = -H_y$, this term is equal to zero [9]. This means that the radiation damping of a free electron accelerated by the plane wave is negligibly small. However, fields in plasmas are distinct from the plane wave. At the laser plasma interface a large longitudinal electrostatic field comparable to the incident laser field can be generated when $n_e \cong n_c/a_0$ where n_e is the plasma density and n_c is the critical density [14]. In addition, in overdense plasma, the laser field is a superposition of the incident and reflected part of the pulse. Different groups of energetic electrons can interact with a different part of the laser field. The energy of an electron accelerated by the plasma electric field and propagating counter to the incident part of the laser pulse can be damped by the incident laser field. The energy of an electron accelerated by the $\boldsymbol{E}\times\boldsymbol{H}$ force or by the wakefield can be damped by the reflected laser field.

The approximation given by Eq. (2) is valid for $e^2a_0\gamma\omega/mc^3 \ll 1$ [9], while the damping force can be the dominant part of the Dirac-Lorentz force, i.e., $e^2a_0\gamma^2\omega/mc^3 > 1$. These conditions are compatible if $\gamma > 1$. However, the parameters of the present simulation are chosen so that $e^2a_0\omega\gamma^2/mc^3 < 1$ and the short wavelength radiation of a particle, which is not resolved by the grid radiation field, still does not affect the other particle motion.

We make a full relativistic 1D particle-in-cell simulation [15,16] with mobile ions for a Cu slab irradiated by a normally incident laser pulse including plasma ionization [15] and the radiation damping in the form of Dirac-Lorentz equation. The laser intensity is varied from 10^{22}–10^{23} W/cm^2, $\lambda = 0.8$ μm, and the pulse duration is 20 fs. The prepulse effect is included through the initial condition for the plasma density. The radiation damping of Eq. (2) is calculated as a friction force self-consistently with the electromagnetic fields. The thickness of the plasma slab is set at 3 μm, as see in Fig. 1, with the maximal ion density $N_i = 9\times10^{21}$ cm^{-3}. The density profile of the plasma corona is an exponential profile with gradient $L = 1.7$ and 0.17 μm indicated by labels 1 and 2 in Fig. 1, respectively. The initial ion charge is $z = 3$ and the initial electron temperature is 100 eV. The time step is chosen to provide the energy conservation better than 0.1% of total laser energy in nonradiating plasmas.

Absorption rates of the Cu plasma are shown in Figs. 2a and 2b. The total absorption rate in the figure is the total pulse energy minus the current electromagnetic pulse energy normalized to the total pulse energy. It is high at about 50%–70%. As seen in Fig. 2b, the absorption rate strongly depends on the density gradient. We attribute

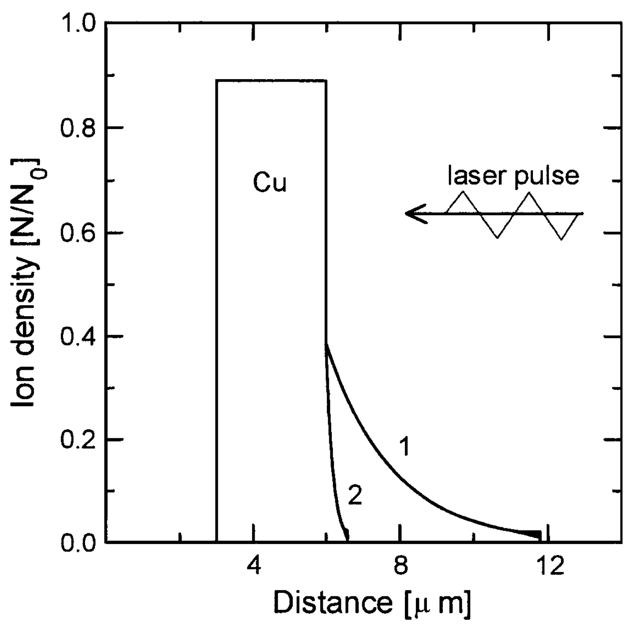

FIG. 1. Initial density distribution in the Cu plasma slab. $N_0 = 10^{22}$ cm^{-3}; the preplasma density $N(x) = 0.4N_0\exp[-(x - x_0)/L]$; (1) $L = 1.7$ μm; (2) $L = 0.17$ μm.

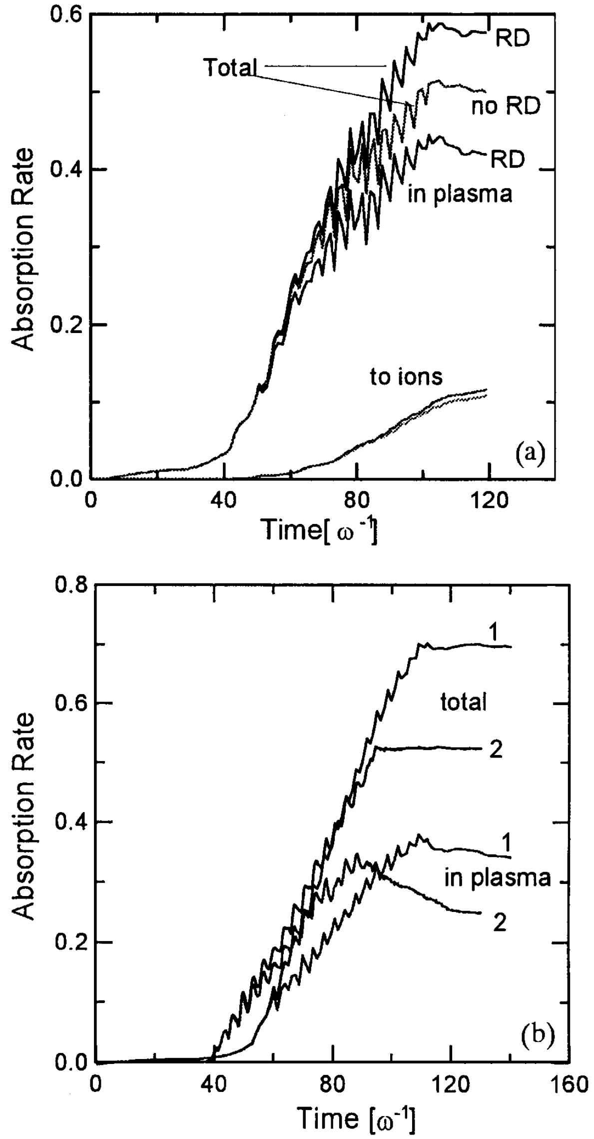

FIG. 2. Temporal evolution of the absorption rate. (a) With (RD) and without (no RD) radiation damping effect, the laser intensity $I = 5 \times 10^{22}$ W/cm^2, the plasma corona gradient $L = 1.7$ μm. (b) With radiation damping effect for the laser intensity $I = 10^{23}$ W/cm^2; (1) $L = 1.7$; (2) $L = 0.17$ μm.

the high absorption rate to the relativistic resonance absorption appearing as a result of the action of the 2ω component of the $\boldsymbol{E} \times \boldsymbol{H}$ force [17]. At the laser intensity $I < 10^{22}$ W/cm^2 we observe no radiation damping effect during the laser irradiation. At laser intensity, $I = 10^{22}$ W/cm^2, the plasma energy loss to radiation is only 0.2%. The effect of radiation becomes strong already at the laser intensity $I = 5 \times 10^{22}$ W/cm^2 ($I\lambda^2 = 3.2 \times 10^{22}$ W μm^2/cm^2). Approximately 55% of the laser energy is absorbed with 40% of the laser energy remaining in the plasma in the form of kinetic and electrostatic field energy after reflection of the laser pulse (indicated by "in plasma" in the figure). The total absorption rate is 7% higher in the presence of RD than that without. This is the pure effect of radiation damping. Ions take 10% of the laser energy. The radiation damping seems to not affect this efficiency. During the laser irradiation, hot electrons with relatively low energy whose motion is not affected by the radiation damping drive the ion acceleration. The effect of radiation damping considerably increases with the laser intensity as shown in Fig. 2b. At the laser intensity $I = 10^{23}$ W/cm^2, from the difference between the total absorbed energy and the energy absorbed in the plasma we see that 35% of the absorbed laser energy is converted to plasma radiation (indicated by 1 in the figure). The pure effect of the radiation damping is 20%. Since this plasma radiation is emitted during the laser pulse interaction, this type of interaction makes possible the generation of a short intense burst of x rays on the order of the pulse duration. The growth of the radiation damping (calculated for $I/10^{22}$ W/cm$^2 = 1, 5, 10$) appears to be proportional to I^2 in contrast to I^3 as expected from the largest radiation damping term in Eq. (3). Setting the force, f, in the energy balance equation, $dQ/dt = -N_{eh} f v$, where N_{eh} is the hot electron density and v is the velocity of the electrons, to be equal to the largest term of Eq. (3), $f_1 \sim r_e^2 c \gamma^2 E^2$ with r_e the electron radius and $\gamma = a_0^2/2$, we expect an energy loss proportional to $Q \sim a_0^6 \sim I^3$ for a constant density of hot electrons N_{eh}. However, the laser radiation cannot penetrate deeply into the overdense plasma and the energy of electrons accelerated by the $E \times H$ force is smaller than the maximal $\gamma = a_0^2/2$ that results in the energy loss dependence on the laser intensity. However, the energy of the electron is high enough for rapid radiation. The radiation time for a relativistic electron can be estimated in a similar way, $d\gamma/dt \sim (e\omega a_0 \gamma)^2/mc^3$. The electron can lose half its energy for radiation in $\Delta t \sim mc^3/(e\omega a_0)\gamma$, for $\gamma = a_0^2/2$ and $a_0 = 100\Delta t$ is 0.2 fs. For the energy acquired this time is less than a few femtoseconds.

The instantaneous spatial distribution of the longitudinal component of the electron momentum in the plasma is given in Fig. 3 without (a) and with (b) RD at $\omega t = 100$. Strong RD cuts the maximal energy of plasma electrons as seen in circles in Fig. 3a. In the bulk plasma, bunches of energetic electrons are propagating into the plasma. The bulk width is determined by the plasma frequency $\omega_{\rm pl}$ which depends on the plasma ionization. In this case, ion charge in the bulk is $z = 20$ (while in the corona the plasma is fully ionized, $z = 29$, due to the optical field ionization) and $\lambda_{\rm pl} = 2\pi c/\omega_{\rm pl} \sim 0.1$ μm. In the plasma corona, electrons are accelerated by the superposition of the laser fields and the plasma electric field. There are three groups of accelerated electrons. The first is the electrons accelerated by the $E \times H$ force of the incoming part of the laser pulse. These electrons propagate into the plasma and can radiate interacting with the plasma field and the outgoing part of the laser pulse. The energy of these electrons exceeds $\gamma = 400$. This means that $f_{\rm RD}/f_{\rm L} = 2e^2\omega\gamma^2 a_0/(3mc^3) \sim 0.25$. As a result, the radiation damping leads to a considerable decrease of the energy of the high energy electrons. The second group

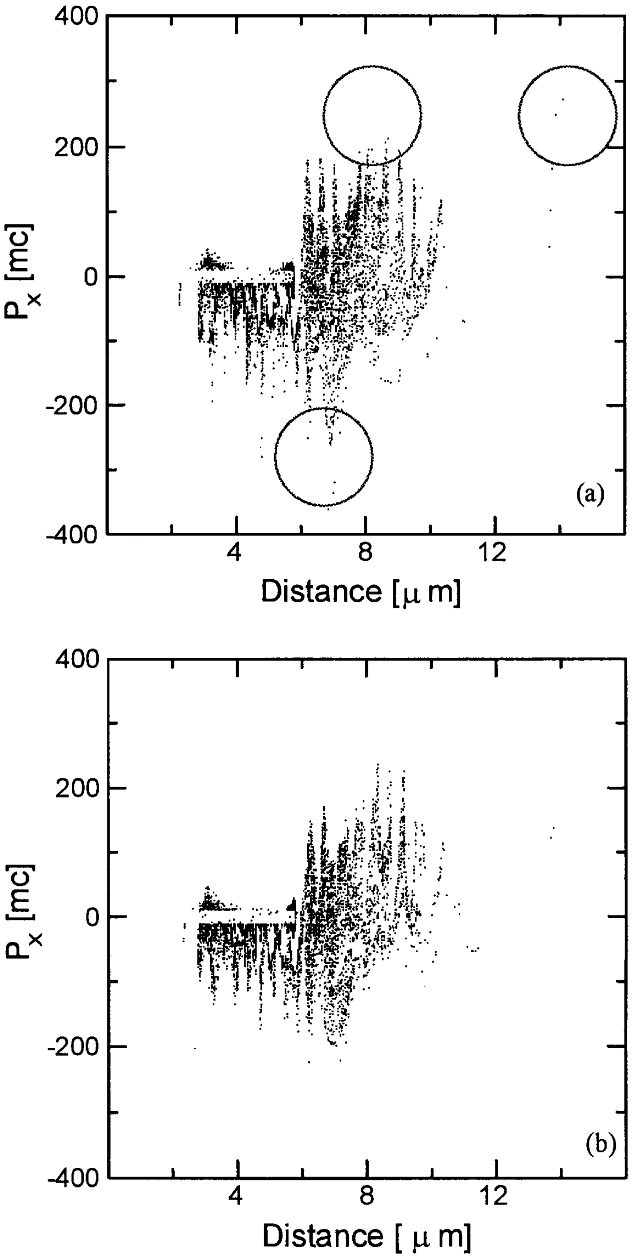

FIG. 3. Spatial distribution of electron longitude momentum P_x at $\omega t = 100$ without (a) and in the presence (b) of the radiation damping effect at the laser intensity $I = 5 \times 10^{22}$ W/cm^2.

is the electrons accelerated by the outgoing part of the laser pulse. Since the reflected laser field is weaker than the incident, the effect of radiation damping is smaller. The third group is the electrons accelerated by the plasma field. The spatial evolution of the plasma field shows, as expected, that the longitudinal plasma field compensating the charge separation produced by the $\boldsymbol{E} \times \boldsymbol{H}$ force is as strong as the laser field approaching $a_0 \sim 80$. Electrons from the return current accelerated by this field can radiate efficiently interacting with the laser field.

In conclusion, we have observed the strong effect of radiation damping on the interaction of an ultrarelativistic laser pulse with overdense plasmas via a fully relativistic particle-in-cell simulation. We have found that the radiation efficiency increases nonlinearly and appears to increase as I^2 in the range of laser intensities $I = 10^{22}$–10^{23} W/cm^2. More than 35% of the laser energy can be converted to a x-ray burst propagating both inward to the plasma and outward at a laser irradiance of $I\lambda^2 = 5 \times 10^{22}$ W μm^2/cm^2 during the laser irradiation. Such a short pulse intense x-ray burst source can have many uses. In addition, these conditions are very close to the classical electrodynamics limit for radiation damping. Therefore, sufficiently high intensity lasers interacting with high-density plasmas may enable us to study the breakdown of classical electrodynamics under laboratory conditions. In future work, more detailed study of the scaling of the absorption efficiency as well as the effects of radiation damping in higher dimensions will be done.

We thank Dr. Y. Ueshima for fruitful discussion.

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