Lasetron: A Proposed Source of Powerful Nuclear-Time-Scale Electromagnetic Bursts

A. E. Kaplan
Electrical and Computer Engineering Department, The Johns Hopkins University, Baltimore, Maryland 21218

P. L. Shkolnikov
Electrical and Computer Engineering Department, SUNY at Stony Brook, Stony Brook, New York 11794
(Received 16 May 2001; published 30 January 2002)

Electromagnetic bursts of substantial energy on a nuclear time scale of $10^{-21}-10^{-22}$ s [zeptosecond (zs) to sub-zs] can theoretically be generated by a perawatt or multiterawatt laser beam focused on a subwavelength-size solid particle or thin wire. Terawatt laser in a similar setup could be instrumental in reaching the subattosecond domain. The system may also generate a half-cycle pulse magnetic field on astrophysical scale up to $\sim 10^6$ T.

DOI: 10.1103/PhysRevLett.88.074801 PACS numbers: 41.60.–m, 41.75.Jv, 42.62.–b, 42.65.Re

Ever shorter electromagnetic (EM) pulses have always been of keen interest, largely as a means of investigating and controlling ever faster processes. Recent proposals [1] explored various avenues to attaining the shortest, subfemtosecond ($10^{-16}-10^{-17}$ s), EM pulses of atomic time-scale duration. In the most recent breakthrough work [2], the train of ~ 0.25 fs pulses have been observed experimentally. The further scale of fundamental interest is that of strong nuclear interactions. Estimating the time scale of a process as $\hbar/(\text{energy scale})$, and keeping in mind that atomic energy scale, defined as the atomic ionization limit, is $\sim 10-20$ eV, while the nuclear energies are beyond 1 MeV, one finds that the nuclear time scale is shorter by about 5 orders of magnitude, i.e., in the $10^{-21}-10^{-22}$ s domain [zeptosecond (zs) to sub-zs]. The feasibility of such pulses has not been considered yet.

In this Letter, we demonstrate theoretically that zs and sub-zs pulses can in fact be generated using petawatt lasers, while already available terawatt lasers may generate subattosecond pulses of $\sim 10^{-19}$ s. The pulses will be radiated by ultrarelativistic electrons driven by circularly polarized high-intensity laser fields. They are basically reminiscent to synchrotron radiation; no synchrotron, however, can even come close to running electrons with the energy of 50 MeV at the (laser) frequency $\omega_L \sim 10^{15}-10^{16}$ s^{-1} in the 0.1 μm radius orbit, as a petawatt laser can. The major distinct feature here is the forced synchronization of all radiating electrons by the driving laser field. Radiation of such a synchronized bunch would be viewed by an observer at any point in the rotation plane as huge pulses/bursts of EM field as short as

$$\tau_{\text{pl}} \sim 1/(2\omega_L \gamma^3), \tag{1}$$

where γ is the electron's relativistic factor. With $\lambda_L \equiv 2\pi c/\omega_L \sim 1$ μm and $\gamma \sim 64$ (attainable with a petawatt laser), we have $\tau_{\text{pl}} \approx 10^{-21}$ s. We call such a system "lasetron." It can be achieved by placing a solid particle or a piece of wire of subwavelength cross section in the focal plane of a superpowerful laser. In addition to zeptosecond pulses with substantial energy, the magnetic field at the center of rotation may reach $\sim 10^6$ T—comparable to fields in the vicinity of white dwarves. Our results also show that the coherent radiation friction drastically limits the rotation energy of electrons in ultraintense laser fields.

Relativistically intense laser field interactions with free electrons have been extensively explored (see, e.g., Ref. [3] and references therein). Closer to our subject, Ref. [4] suggests using electrons in high-intensity, circularly polarized laser light to generate high-order harmonics. However, coherent generation of high-intensity ultrashort pulses by free electrons in laser fields, especially when the radiated power is so high that the radiation damping would dramatically affect the electron motion, has not, to our knowledge, been addressed thus far.

It is known (see, e.g., [5]) that in a circularly polarized EM pulse an electron moves in a helix. We neglect the longitudinal motion by considering only a quasi-steady-state situation when the electron moves in a circle for the time period much longer than one laser cycle. This simple, while advantageous, configuration might be created, for instance, by two circularly polarized counterpropagating laser pulses [6]. Recent results [7] on a similar configuration with a thin film target predict the formation of a very thin layer of free electrons with the density $n_e \sim 10^{23}$ cm^{-3}.

A circularly polarized laser field with amplitude E_L drives an electron with the speed v in a circle with the radius $\rho = \text{v}\lambda_L/2\pi c$ in phase with the field, so that its momentum $p = \text{v}\gamma m_e$ and relativistic factor $\gamma \equiv (1-\beta^2)^{-1/2}$, where $\beta = \text{v}/c$, are, respectively, as [8]

$$p = \varepsilon_L m_e c, \quad \gamma = \sqrt{1+\varepsilon_L^2}, \quad \varepsilon_L \equiv E_L/E_{\text{rel}}, \tag{2}$$

where $E_{\text{rel}}(\omega_L) = m_e \omega_L c/e \approx 10^4 \lambda_L^{-1}$ esu is a relativistic scale of the field strength. (We will see below, however, that Eq. (2) should be amended for high radiation losses.) A rotating electron will radiate the power [6,8] P_e,

$$P_e = m_e c^2 \omega_L \Gamma_e \gamma^2(\gamma^2 - 1), \quad \Gamma_e = (4\pi r_e/3\lambda_L), \tag{3}$$

where $r_e = e^2/m_e c^2 \approx 2.82 \times 10^{-13}$ cm is the classical electron radius. The radiation is concentrated in the angle $\theta \sim 1/\gamma$ around the direction of the instantaneous electron velocity, and an observer in the rotation plane may see only ultrashort bursts of radiation of the duration $\tau_{\rm pl}$, Eq. (1), separated by the laser period, $2\pi/\omega_L$. The Fourier spectrum of the bursts spreads up to the (classical) cutoff

$$\omega_{\max} \sim 3\gamma^3 \omega_L. \tag{4}$$

Quantum cutoff frequency, that of the photon with the energy equal to the electron rotation energy, is $\omega_{\rm qm} \sim m_e \times c^2(\gamma - 1)/\hbar$. They coincide at $\gamma \sim (m_e c^2/3\hbar\omega_L)^{1/2} = O(10^3)$; therefore, even for a petawatt laser with $\gamma \sim 100$ (see below), the radiation near $\omega_{\rm qm}$ is negligible.

We estimate the parameters of the lasetron radiation for coherent radiation sources focused tightly to a few-wavelength spot size $w_L = j\lambda_L, j = O(1)$. One has then $\gamma^2 \approx 1 + 4.6 \times 10^{-11} P_L/j^2$, where P_L is the driving laser power. For simplicity, we assume $j^2 \approx 4.6$ (the beam waist area is then $\approx 7\lambda_L^2$), so that

$$\gamma^2 = 1 + 10 P_L \text{ (TW)}. \tag{5}$$

For quantitative estimates, we consider the following model sources: (i) $PL - P_L = 10^{15}\ W$ (petawatt) laser at $\lambda_L = 1\ \mu$m, a close approximation to the LLNL petawatt laser and a similar system under construction in Japan; (ii) MTW (multiterawatt)—a 100 TW CO_2, $\lambda_L \approx 10\ \mu$m system under construction in Japan; the CO_2 laser at 40 TW is in operation; (iii) LTW—quite widespread lasers of few-TW power; as an example, we will use a 5 TW system at $\lambda_L = 0.8\ \mu$m; and (iv) RK—a relativistic klystron under development at NRL, potentially a 1 TW system with $\lambda_L \approx 3$ cm. Table I illustrates lasetron radiation parameters for a single electron for all these sources. Thus, a single electron in the focus of the petawatt laser would radiate a macroscopic power of 180 W in nuclear time-scale bursts, $\tau_{\rm pl} = 0.26$ zs. The classical cutoff, $\hbar\omega_{cl} = 3$ MeV, lies above the energy threshold of some *photonuclear reactions,* e.g., neutron photoproduction on Be (1.7 MeV). This indicates the potential of lasetron for time-resolved photonuclear physics—provided that a burst carries sufficient energy.

Unfortunately, even for PL, the energy radiated by one electron in one burst is still very low. To increase it substantially, many electrons have to radiate coherently. A straightforward solution is to place a sub-λ_L size solid particle in the focal plane of a high-power laser, which would fully ionize the particle within a fraction of a laser cycle. Free electrons will then experience an "orbital sander" rotation, moving in phase with the field in identical but shifted circular orbits, their relative positions fixed. The resulting radiation will be almost fully coherent, with the radiated power scaling as the particle number squared,

$$P_{\rm rad} \approx N_e^2 P_e. \tag{6}$$

Now, however, P_e cannot be taken from Eqs. (2) and (3), because we have to take into account a new factor—*coherent* radiation friction, or back-reaction of radiation, which is a major player in the phenomenon we consider. Indeed, applying Eq. (6), with γ for P_e taken from Eq. (2), for N_e as small as 10^7, one obtains $P_{\rm rad} = 18$ PW—much higher than the full driving power. The reason for this contradiction is that, when a electron cloud radiates coherently, the radiation losses per electron are much larger than those for incoherent radiation. Therefore, for a sufficiently large number of electrons, the coherent radiation friction must be taken into account from the very beginning.

In its general formulation, this problem has not been solved yet (see, e.g., [9]); we will address it in detail elsewhere. To account for the coherent radiation in this Letter, we approximate a small and dense electron cloud in a strong laser field by a *single pointlike particle* with the charge $q = N_e e$ and mass $m = N_e m_e$, which we call a "fat electron." This model is appropriate because of the strictly field-driven nature of the cloud motion, whereby the tight cloud can be ascribed a single trajectory identical to trajectories of individuals electrons. As discussed above, we neglect the electron motion along the direction of laser propagation; the equation of motion in the plane normal to that axis is then

$$d\vec{p}/dt + \Gamma_{\rm fat}\omega_L \vec{p} \approx e\vec{E}_L, \qquad \Gamma_{\rm fat} \equiv N_e \Gamma_e \gamma^3, \tag{7}$$

where $\Gamma_{\rm fat}$ and $-\Gamma_{\rm fat}\omega_L\vec{p}$ are the radiation damping constant and the "radiation friction" force of the fat electron, respectively (for a single electron, see, e.g., [6]). At the equilibrium between the friction and driving forces, the relativistic factor, instead of (2), becomes now

$$\gamma = \sqrt{1 + \varepsilon_L^2/(1 + \Gamma_{\rm fat}^2)}. \tag{8}$$

For a *single* electron, we have $\Gamma_{\rm fat} \ll 1$, and Eq. (2) is accurate for laser power up to thousands of petawatts. As N_e increases, however, $\Gamma_{\rm fat}$ grows proportionally. One can understand this also in terms of the ratio of energy radiated by a fat electron per radian of rotation angle, $W_{\rm rad} = N_e^2 \Gamma_e m_e c^2 \gamma^2(\gamma^2 - 1)$, to its kinetic energy $W_{\rm rot} = N_e m_e c^2(\gamma - 1)$:

$$\Gamma_{\rm fat} = (1 - \gamma^{-1})^{-1}(W_{\rm rad}/W_{\rm rot}), \tag{9}$$

hence, the increase in $\Gamma_{\rm fat}$ signifies the increase of radiation energy compared to that of the driven rotation; at $\gamma \gg 1$, $\Gamma_{\rm fat} \approx W_{\rm rad}/W_{\rm rot}$. Introducing an "EM size" of the fat electron as $r_{\rm fat} = r_e N_e \gamma^3$, we see that $\Gamma_{\rm fat} = 1$ when $r_{\rm fat} = (3/4\pi)\lambda_L$. The conventional theory of electron radiation at $\hbar\omega \ll m_e c^2$ is based on the fact that

TABLE I. Single-electron output of model sources.

Source	P_e	$\tau_{\rm pl}$	$\hbar\omega_{\max}$	Burst
PL	180 W	0.26 zs	3.9 MeV	0.3 eV
MTW	0.02 W	0.81 as	13 keV	0.1 eV
LTW	0.6 μW	0.6 as	1.7 keV	1.9 μeV
RK	1.5 pW	220 fs	5 meV	0.3 μeV

$r_e \lll \lambda$; i.e., the radiation energy is very small compared to a kinetic one; it breaks down completely if $r_{\rm fat} = O(\lambda)$. In the weak radiation case, a strong driving field, $\varepsilon_L^2 \gg 1$ would result in $\gamma \approx \varepsilon_L$. However, if the radiation is strong and coherent, the further increase of electron energy γ is drastically inhibited as ε_L^2 increases:

$$\gamma \approx (\varepsilon_L/\Gamma_N)^{1/4} \qquad \text{if } \varepsilon_L^2 \gg \Gamma_N^2 + \Gamma_N^{-2/3}, \tag{10}$$

where $\Gamma_N = N_e \Gamma_e$. The strong radiation friction results in a trade-off between the pulse duration and radiated energy: More energy requires more electrons, which in turn limits γ and $\tau_{\rm pl}$. This still allows for spectacular output. For example, if N_e is such that, for $\varepsilon_L = 100$ (PL), Eq. (8) yields $\gamma \approx (2/3)\varepsilon_L$, then $N_e \approx 300$, and one may expect EM bursts of 0.9 zs, separated by 3 fs intervals, each burst carrying 3 fJ energy with the spectral cutoff at 1.2 MeV. If $N_e = 21\,000$, the energy/burst grows to 5 pJ, but γ drops to $(1/4)\varepsilon_L$, so that $\tau_{\rm pl} \sim 17$ zs is still very short.

This trade-off can be circumvented by combining the lasetron with a relativistic heavy ion accelerator. If the lasetron pulses are directed toward a uranium ion beam with $\gamma_{\rm nucl} \approx 100$, as in RHIC at BNL, then uranium nuclei would see in their rest frame the Doppler up-shifted pulses of $\tau_{\rm pl}/2\gamma_{\rm nucl}$ duration. For the second example above, the up-shifted spectrum is concentrated in the area of uranium giant dipole resonance (the maximum of photofission cross section), while the burst duration is short compared with the lifetime of fissioning nuclei. Therefore, a combination of the petawatt laser driven lasetron and RHIC holds a potential for time-resolved measurements and control of fast nuclear fission. Moreover, the nuclei accelerated in the future Large Hadron Collider ($\gamma_{\rm nucl} \approx 3000$) would see lasetron pulses shortened to yet another, yoctosecond (10^{-24} s) domain.

Several mechanisms, which may potentially broaden the lasetron burst or otherwise obscure the proposed effects, require a separate investigation. The phase shift of the driving field along the propagation axis may cause the broadening. In this regard, it is important that two counterpropagating beams create a standing wave, so that the phase of rotation stays almost the same along the beams' axis. The role of Coulomb forces and of the scattering of the rotating electron cloud off the slow moving ion core may also be important. Our preliminary estimates show, however, that it would take up to 10^2 cycles for a cloud to blow up to the radius of its orbit due to Coulomb expansion. The ponderomotive drift of electrons from the laser focus due to gradient force could be fully suppressed by positioning the target precisely in the focus. Even if this is not done, the cloud of electrons would drift as a whole due to the subwavelength size of the electron orbit, thus greatly reducing the unwanted effect due to the drift; the Coulomb field of the core will also prevent the drift. Another factor, a broad-band thermal radiation emitted by the "hot" target plasma may not hamper the observation of the zeptosecond bursts, because of its almost continuum spectrum, whereas the lasetron bursts have the periodicity of the laser cycles and thus a multi-narrow-line spectrum (the relativistic harmonics of laser). The lasetron effect observation can then be enhanced by using the laser-cycle analogy of a boxcar detector; together with a very high collimation of the lasetron radiation (see below), this makes a potential tool for the resolution of lasetron bursts on the thermal background.

A *thin wire* positioned in the laser focal plane normally to the laser beam propagation, Fig. 1, could be an even more promising target. Our calculations show that its ultrashort burst radiation will be coherent. Its full energy at the nth harmonic with the polarization parallel to the plane of electron rotation into solid angle $d\Omega$ is as

$$\begin{aligned} dW_n/d\Omega \approx\ & (27/4\pi^2) m_e c^2 n^{-4/3} n_e^2 \lambda_L r_e h^2 x_n^2 \sin^2(ndk_L/2) \\ & \times [\sin(\zeta)\zeta^{-1}]^2 \exp[-(k_L n x_n \psi)^2] \\ & \times \exp\{-[(n^{2/3}+1)(\gamma_{\rm max}^{-2}+\theta^2)/2]^2\}. \end{aligned} \tag{11}$$

Here, $k_L = 2\pi/\lambda_L$ is the laser wave number; n_e is the electron density; h and d are the thickness (along the axis of laser propagation) and the width (in the cross section of laser beam) of the wire, respectively (we assume $h, d \ll \lambda_L$); θ is the angle between the direction of observation and the plane of electron rotation; $\zeta = (nhk_L/2)(\theta - 1/n)$ is angle θ normalized and centered for the nth harmonics; ψ is the angle between two planes: the one formed by the axis of laser beam and the direction of observations, and the other normal to the wire; and x_n is the effective "antenna size" for each number n and angle θ. If $n, \varepsilon_{\rm max}^2 \gg 1$, we have $x_n \approx w_L\{\ln/\varepsilon_{\rm max}^2(2n^{-2/3} - \theta^2)]\}^{1/2}$; and $\gamma_{\rm max}$ is the relativistic factor at the field maximum $\varepsilon_{\rm max}^2$. The function $\gamma_{\rm max}$ ($\varepsilon_{\rm max}$) is affected by radiation friction to a lesser degree than in the fat electron effect in a compact cloud; it will be considered by us in detail elsewhere.

Because of the coherence, the wire antenna will radiate only twice a laser cycle, with the radiation highly concentrated in *two very narrow beams* strictly normal to the wire and almost normal to the laser beam (see Fig. 1). (This is particularly beneficial for the above-described combination of the lasetron and an ion accelerator.) In the plane normal to the wire, these beams are slightly tilted toward the propagation of the laser beam. The angular collimation of the radiation by such a 3D antenna due to the laser beam of the size $w_L \gg \rho$ would be much greater than that of a single electron or compact cloud; it will mostly be concentrated within the angle,

$$\Delta\psi_{\rm min} \sim (\lambda_L/2\pi w_L)\varepsilon_{\rm max}^{-3}. \tag{12}$$

This high collimation of radiation will result in great enhancement of radiation intensity in the far field area. The pulses appearing only in two well-defined opposite directions, and separated in time by half the laser cycle, would also provide a clear signature of the lasetron effect.

To evaluate the *magnetic (M) field* in a lasetron, we consider a spherical target of a small diameter $d \ll \lambda_L$, placed in the tightly focused laser field. The driven motion of the ionized electron cloud, which will largely maintain its initial small size for a large number of laser cycles,

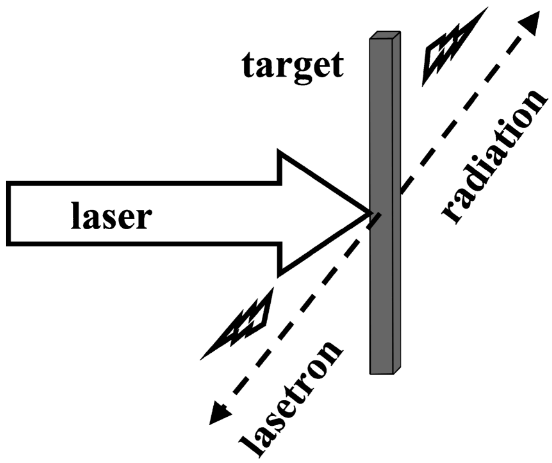

FIG. 1. Geometry of the wire target lasetron radiation. Ultrashort bursts propagate along the dashed lines normal to the laser propagation and to the target (wire).

will create a strong M field normal to the rotation plane. At the center of rotation, this field will have a quasistationary component, to the extent following the laser amplitude envelope, with the M-field induction in Gauss as

$$B \approx eN_e\beta\rho^{-2} = (3/2)B_{\rm rel}\beta^{-1}N_e\Gamma_e\,, \tag{13}$$

where $\rho = \beta k_L^{-1}$ is the radius of the cloud orbit, and $B_{\rm rel} \equiv mc\omega_L/e$ is a relativistic scale of M field, equal to $E_{\rm rel}$ [Eq. (2)]. Assuming the upper limit of the cloud size to be $d \sim \rho$, and, hence, the number of electrons $N_e \sim \pi n_e \rho^3/6$, and the excitation to be sufficiently relativistic, i.e., $\beta \sim 1$ (see also below), and $\rho \sim k_L^{-1}$, we estimate the highest possible M field in the lasetron as

$$B_{\rm max} \sim en_e\lambda_L/12 = \pi/6 \cdot (\lambda_L/\lambda_C)(n_e a_0^3) \cdot B_0\,, \tag{14}$$

where $B_0 = e\alpha/a_0^2 \sim 1.33 \times 10^5$ G is the "Bohr" M-field scale calculated as a M field induced by a *classical* electron rotating with the speed $\beta = \alpha$ (corresponding to the classical speed of an electron at the ground state of an H atom) and Bohr radius $a_0 = 0.53$ Å, $\alpha \approx 1/137$ is the fine structure constant, and $\lambda_C \approx 2.4 \times 10^{-10}$ cm is the Compton wavelength. In Eq. (14), we have $\lambda_L/\lambda_C \ggg 1$ and $n_e a_0^3 = O(1)$, so that $B_{\rm max} \ggg B_0$. Choosing a high-Z electron-righ material to be highly ionized in a superstrong laser field, we can assume $n_e \sim 10^{24}$ cm^{-3}, so that $B_{\rm max} \sim 4 \times 10^9$ G for $\lambda_L \sim 1$ μm, and $\sim 4 \times 10^{10}$ G for $\lambda_L \sim 10$ μm. The field is parallel to the laser propagation direction, and has the transverse size $\sim 2\rho \sim \lambda_L/\pi$; thus the total magnetic flux is $2\pi\rho^2 B_{\rm max} \propto \lambda_L^3$. The M-field duration is about the same as that of the originating laser pulse.

Note that, for generating high M fields, there is no need for too high γ; the condition $\beta \sim 1$ is satisfied if $\varepsilon_L^2 > \varepsilon_{\rm rel}^2 \equiv 1 + N_e^2\Gamma_e^2$, while $\varepsilon_L^2 \gg \varepsilon_{\rm rel}^2$ is required for generating ultrashort pulses. Even in a nonrelativistic situation, $\varepsilon_L < \varepsilon_{\rm rel}$, the M field can be much higher than that for any available Earth sources. In this case, the right-hand side in Eq. (14) must be corrected by the factor $\beta^2 \sim \varepsilon_L^2/\varepsilon_{\rm rel}^2$, which for $\beta \sim 0.16$ (electron energy $\sim$6 keV) still results in $B \sim 10^8$–10^9 G. The lasetron M field can be observed by scattering of neutron beams by the cloud. Another way of its detection could be the observation of a magnetic single-cycle pulse excited by the half-cycle pulse of magnetic dipole $B(t)$, to be discussed by us elsewhere, similarly to a pulse excited by the half-cycle pulse of an elementary electric dipole. In both cases, the temporal profile of the pulse mimics the time derivative of the originating dipole pulse [10].

In conclusion, we have demonstrated theoretically the feasibility of a system (lasetron) capable of generating EM bursts of large energy on a nuclear time scale 10^{-21}–10^{-22} s using available lasers. It is also capable of generating a superstrong magnetic pulse field on astrophysical scale up to $\sim 10^{10}$ G.

This work is supported by AFOSR.

[1] A. E. Kaplan, Phys. Rev. Lett. **73**, 1243 (1994); P. B. Corkum, N. H. Burnett, and M. Y. Ivanov, Opt. Lett. **19**, 1870 (1994); A. E. Kaplan and P. L. Shkolnikov, Phys. Rev. Lett. **75**, 2316 (1995); A. E. Kaplan, S. F. Straub, and P. L. Shkolnikov, J. Opt. Soc. Am. B **14**, 3013 (1997); S. E. Harris and A. V. Sokolov, Phys. Rev. Lett. **81**, 2894 (1998); I. P. Christov, M. M. Murnane, and H. C. Kapteyn, Opt. Commun. **148**, 75 (1998); E. V. Kazantseva and A. I. Maimistov, Phys. Lett. A **263**, 434 (1999); A. V. Sokolov, D. D. Yavuz, and S. E. Harris, Opt. Lett. **24**, 557 (1999); F. Le Kien *et al.,* Phys. Rev. A **60**, 1562 (1999); N. A. Papadogiannis *et al.,* Phys. Rev. Lett. **83**, 4289 (1999).

[2] P. M. Paul *et al.,* Science **292**, 1689 (2001); E. Hertz *et al.,* Phys. Rev. A **64**, 051801 (2001).

[3] F. V. Hartemann, Phys. Plasmas **5**, 2037 (1998).

[4] W. Yu *et al.,* in *X-ray Lasers—1996,* IOP Conf. Proc. No. 151 (Institute of Physics and Physical Society, London, 1996), p. 460.

[5] E. S. Sarachik and G. T. Shappert, Phys. Rev. D **1**, 2738 (1970).

[6] A. E. Kaplan, Phys. Rev. Lett. **56**, 456 (1986); A. E. Kaplan and Y. J. Ding, IEEE J. Quantum. Electron. **24**, 1470 (1988).

[7] B. F. Shen and J. Meyer-ter-Vehn, Phys. Plasmas **8**, 1003 (2001).

[8] L. Landau and E. Lifshitz, *Classical Field Theory* (Pergamon, New York, 1975); J. D. Jackson, *Classical Electrodynamics* (Wiley, New York, 1975).

[9] I. Kimel and L. R. Elias, Phys. Rev. Lett. **75**, 4210 (1995); Nucl. Instrum. Methods Phys. Res., Sect. A **375**, 565 (1996).

[10] A. E. Kaplan, J. Opt. Soc. Am. B **15**, 951 (1998).