Evidence of Ultrashort Electron Bunches in Laser-Plasma Interactions at Relativistic Intensities

S. D. Baton, J. J. Santos, F. Amiranoff, H. Popescu, L. Gremillet,* M. Koenig, E. Martinolli, and O. Guilbaud
Laboratoire pour l'Utilisation des Lasers Intenses, UMR 7605 CNRS-CEA-Ecole Polytechnique-Université Paris VI, Palaiseau, France

C. Rousseaux and M. Rabec Le Gloahec
Commissariat à l'Energie Atomique, Bruyères-le-Châtel, France

T. Hall
Department of Physics, University of Essex, Colchester, United Kingdom

D. Batani, E. Perelli, and F. Scianitti
Dipartimento di Fisica "G. Occhialini" and INFM, Università di Milano-Bicocca, Italy

T. E. Cowan
General Atomics, San Diego, California, USA
(Received 31 October 2002; published 5 September 2003)

The second harmonic of the laser light ($2\omega_0$) is observed on the rear side of thick solid targets irradiated by a laser beam at relativistic intensities. This emission is explained by the acceleration by the laser pulse in front of the target of short bunches of electrons separated by the period (or half the period) of the laser light. When reaching the rear side of the target, these electron bunches emit coherent transition radiation at $2\omega_0$. The observations indicate that, in our conditions, the *minimum* fraction of the laser energy transferred to these electron bunches is of the order of 1%.

DOI: 10.1103/PhysRevLett.91.105001 PACS numbers: 52.38.Kd, 41.75.Jv, 52.70.Kz

The interaction of a high-intensity laser beam with a plasma can lead to the acceleration of electrons to relativistic energies. This may have important applications in various domains such as laser particle acceleration [1], the fast igniter scheme for inertial confinement fusion [2–4], the generation of intense and short duration γ-ray sources for radiography [5], and sources of fast ions [6].

The characteristics of these electron sources depend on the laser and plasma parameters and on the geometry of the interaction. The accelerating field can be (i) the laser electric field itself, with an important effect of the laser magnetic field at high laser intensities, (ii) low frequency electric and magnetic fields, or (iii) high frequency electron plasma waves.

For p-polarized pulses obliquely incident on solid targets with a sharp density gradient, the component of the laser electric field perpendicular to the target surface and the space charge field in the plasma in front of the solid target both play a dominant role. Vacuum heating in different regimes can lead to the injection in the target of electron bunches separated by the laser period [7–11]. In the same geometry, resonant absorption [12] excites a plasma wave at the critical density. The breaking of this wave ejects fast electrons mainly towards the vacuum region. However the separation between the electron bunches is controlled by the wave-breaking time which is not directly correlated with the laser period [13].

At normal incidence on a sharp density gradient and high laser intensities, the magnetic term of the Lorentz force ($ev \times B$) (directed along the target normal) becomes comparable to the electric force (eE) and can play a role comparable to the role played by the electric field in oblique incidence. In this case, the so-called $J \times B$ heating [14,15] accelerates bunches of electrons in the target at twice the laser frequency [16–20]. At very high intensities in longer scale length plasmas or moderate density plasmas (a few times the critical density), hole boring curves the interaction surface. Thus even at normal incidence locally p-polarized light can give rise to vacuum heating with the ejection of electron bunches at the laser frequency perpendicular to the local interaction surface [16,18,21].

In much larger undercritical plasmas, Raman-like instabilities excite electron plasma waves at the plasma frequency. Wave breaking of these waves will accelerate electron bunches separated by the plasma period, much longer than the laser period itself [22].

The bunching of the accelerated electrons has not been observed yet. The identification of vacuum heating at oblique incidence and moderate intensities has been made indirectly by measurement of the absorption coefficient [23]. In this Letter, we present the first evidence of acceleration of electron bunches by vacuum or by $v \times B$ heating in laser solid interaction experiments.

The experiment was performed on the LULI 100 TW laser facility. The 400 fs, 1.057 μm laser pulse with energy up to 20 J was focused by a $f/3$ off-axis parabola at normal incidence onto Al targets. The laser focal spot was 15 to 20 μm in diameter corresponding to a maximum intensity of 3×10^{19} W/cm^2. Because of the 1–2 ns long pedestal preceding the main pulse at an intensity of 10^{12}–10^{13} W/cm^2 (i.e., a contrast in terms of energy of $1{:}10^{-4}$), a 30–50 μm long plasma was formed in front of the solid target [24]. The optical radiation of the rear side of the target was collected on axis with $f/3.2$ optics. The image of the emitting region was sent onto a charge-coupled device camera with a magnification of 24 and into the entrance slit of a spectrometer-S20 streak camera combination with a magnification of 1. The spectral resolution was limited by the slit width to 5 nm and the time resolution was $\approx$ 10 ps. The sensitivity of the entire system was established with an absolutely calibrated emission lamp.

In preceding experiments, time-integrated visible images of the rear side of the target revealed two features [25]. A spatially diffuse signal is due to late time (~ns after the laser pulse) emission of the cooling and expanding plasma. In the center of this broad emission there was a bright localized signal lasting only a few picoseconds [25], which was explained as either (i) optical transition radiation (OTR) emitted by electrons leaving the target or (ii) bremsstrahlung radiation emitted by the electrons when they circulate outside the target before being pulled back into the target by strong electric fields. In these experiments, as a precaution against stray light at the second harmonic of the laser frequency, the spectral region around 0.53 mm was excluded from the detectors using optical color filters.

In this Letter we analyze the spectrum of the central bright spot but with the filtering around 0.53 mm removed. A time-resolved spectrum is shown in Fig. 1. The radiation consists of a wide spectrum and an intense and narrow contribution near 530 nm, the second harmonic of the laser light ($2\omega_0$). We made stringent tests to ensure that $2\omega_0$ light really comes from the back side of the foil. For example, we constructed a "black opaque box" around the target, such that scattered light coming from the front of the target was delayed by $>$ 1 ns. Light from the front side coming through the target after shock breakout would similarly be delayed by $\sim$1 ns. The $2\omega_0$ signal was observed to be synchronized with the laser pulse, thus ruling out the possibility that it was from self-emission or harmonic generation in the front side laser-produced plasma.

The integrated energy in the $2\omega_0$ line (after subtraction of the broad continuum) is plotted as a function of Al target thickness in Fig. 2. It decreases rapidly with target thickness but can still be observed after 900 μm of Al. As in [25] the broad spectral signal can be attributed to OTR radiation from the bulk of the electrons passing through the target but, as detailed below, the $2\omega_0$ line is attributed

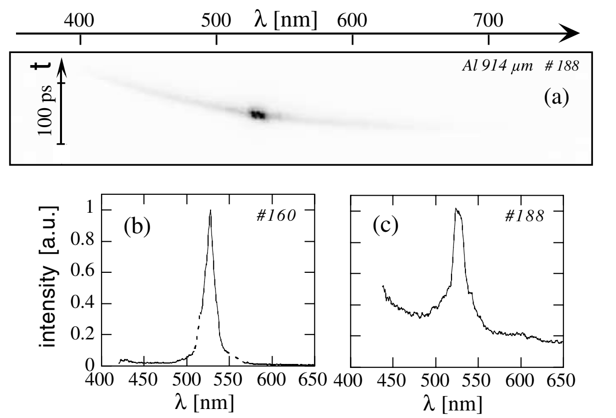

FIG. 1. (a) Image of the spectrum measured on the streak camera for a shot on 914 μm of Al. The apparent distortion in time of this large spectrum is attributed to the dispersion due to the optical components and the streak camera imaging. Two examples of profiles of time-integrated spectra obtained, respectively, for 450 μm (b), and 914 μm (c) Al targets.

to coherent transition radiation (CTR) generated by a relatively small number of high-energy electrons in ultrashort bunches.

The analysis of the data in terms of CTR is detailed in this section. When an electron propagates in a medium, it excites the motion of the background charges, which in turn can radiate. When the electron crosses an interface between two media with different dielectric properties, as in the case of a metal-vacuum transition, the so-called transition radiation emitted near the surface of the material can propagate into the vacuum. In the visible, it is termed OTR [26]. In the case of many independent incident electrons, the total electromagnetic field on a detector outside the material is the sum of the individual fields due to each traveling electron. When these electrons are bunched inside a very short and very narrow bunch, these fields add coherently for wavelengths much larger than the bunch length. This CTR [27–29] leads to a much higher emitted power, roughly proportional to the square of the number of electrons N^2, instead of N in the OTR incoherent case. In addition, when the electrons are

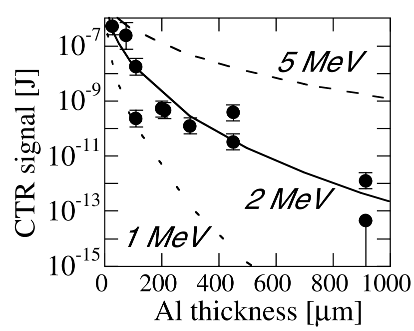

FIG. 2. Integrated energy in the $2\omega_0$ line as a function of Al target thickness. The different curves (for three electron temperatures) correspond to the theoretical estimates explained in the text.

bunched in successive short bunches separated by a time delay δT, the emission is coherent for wavelengths close to $c\delta T$.

In the case of M identical electron bunches separated by a time delay δT, each containing P electrons, and injected perpendicularly to the target surface, the general formula for the coherent intensity collected at the back side of the target at frequency ω is given by

$$I(\omega) = \eta(\omega)P^2|i(\omega)|^2 \frac{\sin^2(M\omega\delta T/2)}{\sin^2(\omega\delta T/2)}. \quad (1)$$

The first term on the right-hand side of $\eta(\omega)$ is the light intensity emitted by a single electron inside the aperture of the collecting optics [in this simple formula, $\eta(\omega)$ is assumed independent of the electron energy, but the actual dependence with the energy and the angular distribution have been taken into account in the final calculation]. The third term is the modulus squared of the Fourier transform of the electron current of a single bunch at the back of the target $i(t)$ (the exact phase of the electric field emitted by each electron is thus taken into account). $i(t)$ is the electron flux (number of electrons per second normalized so that the time integral is equal to 1). The last term corresponds to the coherent addition of the fields generated by each of the M bunches. Equation (1) predicts intense emission lines at the harmonics of the emission frequency (corresponding to a period δT) of the electron bunches. It must be stressed that the spatial distribution of the current on the target surface has been neglected in this formula, as if all the electrons reached the surface at the same point. Taking into account this spatial distribution may have a dramatic effect on the absolute signal. A more precise formula would necessarily take into account the variation of $\eta(\omega)$ with the electron energy, as well as the exact angular spectrum of the electrons and angular pattern of the emission, the energy loss and the angular scattering of the electrons in the material, and the spatial distribution of the electrons at the rear side of the target.

In the theoretical estimates we will further assume that each bunch has an initial duration τ [i.e., initial current $\sim \exp(-t^2/\tau^2)$] and a 1D relativistic Maxwellian velocity distribution. Let $j(t)$ be the current at the rear surface of the target corresponding to an initial current given by a δ function. $j(t)$ is normalized such that $\int j(t)dt = 1$. The third term in Eq. (1) then becomes $|i(\omega)|^2 = |j(\omega)|^2 e^{-\omega^2\tau^2/2}$, where $j(\omega)$ is the Fourier transform of $j(t)$. We will also assume that each electron keeps its initial velocity throughout the whole target thickness. Thus, the main effect is the expansion in the target of each bunch composed of electrons with different velocities as shown in Fig. 3.

The main variation of the intensity with target thickness comes from the Fourier component, which decreases very rapidly as the rise time of $j(t)$ increases. In fact, and except for very small thickness, for which the duration of the signal is shorter than the period of the measured radiation, only the initial part of the current really contributes to the final signal. It means that only the highest energy electrons contribute to the signal, typically those with an energy larger than the temperature (in units of energy) of the 1D Maxwellian distribution.

The incoherent signal on the contrary, is obtained for a random distribution of the electrons over a time much longer than the period of the measured pulsation ω. It is given by a sum in intensity of the signal emitted by each electron:

$$I(\omega) = \eta(\omega)PM = \eta(\omega)N.$$

The existence of the $2\omega_0$ line at the rear of the target is a clear signature of the acceleration of electron bunches at ω_0 or $2\omega_0$ in the laser-plasma interaction region. This is thus the first evidence of the acceleration of ultrashort relativistic electron bunches in laser-plasma interactions.

From the width of the $2\omega_0$ line (~ 20 nm FWHM) and from the absolute signal, information can be obtained on the total energy in the electrons at the origin of the coherent emission. Using the simplest model described above, that of ballistic propagation of the electrons in the target, we calculated the expected signal for different initial electron parameters. In these calculations we took into account the dependence of the intensity $\eta(\omega)$ with electron energy [25]. We will first consider the case of electron bunches emitted at $2\omega_0$ by $v \times B$ heating. If we assume that the width of the line is related to the number of bunches M, a good fit is obtained for $M \approx 30$ which corresponds to a total time duration of ≈ 50 fs. We also assume, as observed in simulations [20], that each bunch lasts only a small fraction of the ω_0 period (we

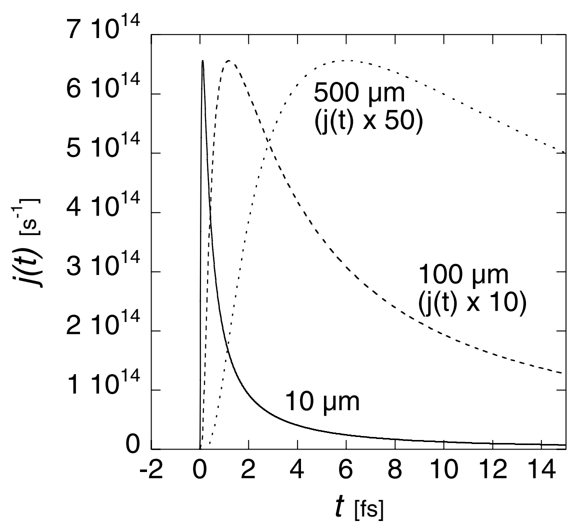

FIG. 3. Evolution of normalized current $j(t)$ for 10 μm target thickness and a 2 MeV electron temperature (solid line). The dashed and dotted lines correspond, respectively, to the evolution of $j(t)$ for 100 μm (multiplied by a factor of 10) and 500 μm (multiplied by a factor of 50).

chose $\tau = \omega_0^{-1}/10 = 0.35$ fs). The two remaining and most important parameters are (i) the number of electrons (or the energy) per bunch, which modifies the amplitude of the signal, and (ii) the temperature of the 1D relativistic distribution, which determines the variation of the signal with target thickness. As can be seen in Fig. 2, a rather good fit is obtained with a temperature of 2 MeV. With this temperature, the line in Fig. 2 has been plotted for 38 μJ per bunch and 30 bunches corresponding to an energy of the order of 4×10^{-4} of the laser energy interacting with the target in half a laser period. The total energy in the $M \approx 30$ bunches which contribute to the coherent signal is $\sim$1.1 mJ, i.e., about 10^{-4} of the total laser energy.

As noted above, even if the targets are irradiated at normal incidence, electron bunches can be emitted at ω_0 by vacuum heating. In this case, using Eq. (1) and changing δT by $\delta T/2$, M, and P, it is easy to show that the same coherent signal at $2\omega_0$ would be obtained for typically 15 bunches of the same duration and 76 μJ per bunch. However, it is reasonable to assume that at the beginning of the pulse, when the laser interacts at normal incidence, part of the electrons are emitted at $2\omega_0$.

Several effects will decrease the intensity of this coherent signal. The most important of them is probably the phase variations due to the effective spatial and temporal distribution of the electrons at the back of the foil, due either to the initial angular divergence of the beam or to angular deviations and energy losses in the foil. CTR could thus be an interesting diagnostic of the electron transport in the target. A more precise calculation taking into account all these effects as well as the angle and time dependent energy spectrum [18–20] and the electric and magnetic fields inside the target (e.g., [30]) is beyond the scope of this paper. It should be noted that, as in [25], part of the $2\omega_0$ signal could be due to synchrotron radiation of the outgoing electrons. However, this does not modify the conclusion on the existence of very short electron bunches.

As discussed in [25], the incoherent part of the spectrum can be explained by about 1 to 2 J of unbunched electrons.

In conclusion, the emission at $2\omega_0$ from the back of a thick target irradiated by a laser beam at high laser intensity has been observed. This is the evidence for the acceleration of ultrashort relativistic electron bunches by the $v \times B$ force or vacuum heating in the interaction region. First estimates show that the *minimum* total energy in these electron bunches is a modest fraction ($\approx 10^{-4}$) of the laser energy but more precise calculations could greatly enhance this value. However, the diagnostic is sensitive only to high-energy electrons, typically larger than a few MeV in our conditions, and it does not give any information on the lower energy part of the spectrum. These high-energy electron bunches could be responsible for the jets observed in previous experiments [31,32]. These clearly visible jets can indeed be due to a small number of relativistic electrons with a modest total energy (< 0.01 J). A more detailed study will be necessary to obtain a precise quantitative estimation of the total energy of the electrons accelerated by the $v \times B$ force and by vacuum heating.

The authors would like to thank the technical staff of LULI and D. Gontier, C. Cholet, and Ph. Mounaix. Special thanks to M. H. Key, J. A. Koch, A. J. Mackinnon, R. R. Freeman, R. A. Snavely, C. Andersen, and R. B. Stephens. Part of this work has been supported by the European Contracts No. LULI ACCESS HPRI-1999-CT 00052 and No. TMR ERBFMGE-CT95-0044 and by Grant No. E1127 from Région Ile-de-France. One of us (J. J. S.) was financed by MCT (Portugal) under the Contract No. PRAXIS XXI BD/18108/98.

*Present address: Commissariat à l'Energie Atomique, Bruyères-le-Châtel, France.

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