Ultrafast target charging due to polarization triggered by laser-accelerated electrons

A. V. Brantov,1,2 A. S. Kuratov,1,2 Yu. M. Aliev,1 and V. Yu. Bychenkov1,2

1P. N. Lebedev Physics Institute, Russian Academy of Science, Leninskiy Prospect 53, Moscow 119991, Russia
2Center for Fundamental and Applied Research, Dukhov Research Institute of Electronics (VNIIA), Moscow 127055, Russia

A significant step has been made towards understanding the physics of the transient surface current generated as a result of the fast loss of hot electrons to the vacuum, the proposed mechanism is associated with excitation of the fast current by electric polarization due to transition radiation triggered by ejected electrons. We present a corresponding theoretical model and compare it with two simulation models using the FDTD (finite-difference time-domain) and PIC (particle-in-cell) methods. Distinctive features of the proposed theory are clearly manifested in both of these models.

Strong ultra-fast surface fields and electric currents in the interaction of short intense laser pulses with solid dense targets are currently of great interest in both fundamental and applied science[1][3]. The generation of a strong surface current is closely related to lateral electron transport, electromagnetic surface waves, and hot spot expansion on the target. The last, for example, is important for ion acceleration. The effect of lateral electron transport is considered in the context of the physics of surface guided schemes for fast ignition[11]. Excitation of electromagnetic surface waves is an elegant way to produce THz radiation with well-concentrated energy[7][8], which is an interesting complement to the widely discussed laser generation of THz radiation into free space[9][10].

We especially emphasize the study with an entire series of experiments with laser-triggered transient electric currents along a target surface[11][4][6], to which our paper is closest. Excitation of a lateral transient current has been standardly associated with a loss of plasma neutrality because the laser-accelerated hot electrons escape to the vacuum. This leads to generating a large positive potential, which provides target neutralization via charge redistribution and the antenna-like propagation of an electromagnetic disturbance away from the interaction region[3]. The discharge of laser-irradiated targets was thus assumed as a physical reason in several previous studies. Unfortunately, no theoretical model has yet been developed to support such a scenario. We do not intend to eliminate this forthcoming and instead aim to highlight another mechanism for generating a lateral transient current, which can be very effective.

Here, we present theoretical model based on the idea that a laser-accelerated electron bunch crossing the target–vacuum boundary generates a fast surface field and the corresponding lateral skin current in the form of a polarization wave. Such a traveling-wave skin current naturally appears in the same approach that was used in the theory of transition radiation[11] and is a fundamental effect for a high-conductivity half-bounded medium, for example, a solid dense plasma. We also complement our theory with two simulation models using FDTD (finite-difference time-domain) and PIC (particle-in-cell) methods. Distinctive features of the proposed theory are clearly manifested in both of these models.

The full set of Maxwell equations for the generated electric field in the Fourier space $(\omega, k)$ is reducible to the equations for the normal ($x$) components $E^p_x(x > 0)$ of this field in a vacuum and $E^p_x(x < 0)$ in a medium (irradiated target),

$$\frac{\partial^2 E^p_x}{\partial x^2} - k^2_p x^2 = \frac{4\pi e \partial n_{\omega,k}^e}{\epsilon x} - \frac{4\pi e i \omega}{c^2} n_{\omega,k}^e,$$

where $k \perp x$, $v = (v, 0, 0)$, $k_v = \sqrt{k^2 - \omega^2 / c^2}$, $k_p = \sqrt{k^2 - \omega^2 / c^2}$, and $\epsilon$ is the dielectric susceptibility. For $n^e(x - vt, r_z)$, the well-known solution[11] of Eq. 1, which takes the continuity of the tangential field components (both electric and magnetic) into account, $\omega$ as $\omega + i0$ and $n \equiv n(\omega, k)$ as given by $n_{\omega,k}^e = ne^{i\omega x} / v$, can
be written as

\[
E_x^v = -\frac{4\pi enkv}{v^2k_v^2 + \omega^2} \left( i\omega \left( 1 - \frac{v^2}{c^2} \right) e^{-i\omega t} + \frac{k(\epsilon - 1)(vk_v + i\omega)(1 - \frac{v^2/c^2}{\epsilon})}{(k_v + ek_v)(vk_v + i\omega)} e^{-ik_vx} \right), \tag{2}
\]

\[
E_y^v = -\frac{4\pi enkv}{\epsilon(\epsilon v^2k_v^2 + \omega^2)} \left( i\omega \left( 1 - \frac{v^2}{c^2} \right) e^{-i\omega t} + \frac{k(\epsilon - 1)(vk_v - i\omega)(1 - \epsilon v^2/c^2)}{(k_v + ek_v)(vk_v - i\omega)} e^{ik_vx} \right).
\]

Solution (2) is a sum of the bunch field and the induced field (the respective first and second terms). From the standpoint of physical effects, Eq. (2) describes the generation of transition radiation into free space \[12\] and transition radiation along a surface \[13\]. Both types of radiation are well studied in the far-field approximation, but the transient electromagnetic fields and electrical current have not yet been studied, because the corresponding theory should be based on the near-field theory. The development of such a theory is our main goal. Experiments both on laser-triggered propagating electromagnetic pulses and on electron transport along the conducting target \[14,15\] are urgently requested. Below, we use two-dimensional geometry in the space \((x, z)\) to analyze this problem.

Generation of an electromagnetic surface wave is defined by the pole in Eq. (2), i.e., the dispersion relation \(d_0 = k_v + ek_v = 0\), which has the solution \(k_v = \pm\omega/\epsilon(1 - 1/(2\epsilon))\) for \(|\epsilon| \gg 1\). This contribution gives the expression for the normal component of the vacuum-side surface wave field:

\[
E_x^v = -\int \frac{d\omega}{2\epsilon} i\epsilon n\omega e^{-i\omega(t\pm\frac{z}{c})} e^{-i\pi|\epsilon|}e^{-i\omega t} \tag{3}
\]

We note that surface field amplitude (3) contains a small factor \(1/\sqrt{|\epsilon|}\) and hence disappears as \(|\epsilon| \to \infty\).

In the limit of an ideal conductor \(|\epsilon| \to \infty (E_y^v = 0)\), the normal component \(E_x^v = E_x^v(x = 0)\) of the vacuum-side field at the surface can be written as

\[
E_{x0}^v = E_r + E_d = -4\pi en \left( \frac{v}{v^2k_v} + \frac{1 - v^2/c^2}{v^2k_v - i\omega} \right), \tag{4}
\]

where the term \(E_r\) is responsible for the transition radiation field and the term \(E_d\) represents the dipole-like field of the uniformly moving charge and its mirror image with respect to the surface \[14\]. For a moving point charge, the spatial Fourier component of \(E_r\) is \(E_r(k) = 4\pi en\nu J_0(\epsilon v t)/c\), where \(J_0\) is a Bessel function, which agrees with the result in \[14\]. In the space \((x, z)\), Eq. (4) becomes

\[
E_{x0}^v = \begin{cases} 
-\frac{4\epsilon v^2e^2\gamma^2}{\sqrt{c^2t^2 - z^2}^2 + v^2t^2\gamma^2}, & z^2 < c^2t^2, \\
0, & z^2 > c^2t^2, 
\end{cases} \tag{5}
\]

where \(\lambda\) is the linear charge density and \(\gamma = 1/\sqrt{1 - v^2/c^2}\) is the Lorentz factor of a moving (along \(x\)) filament (along \(y\)) charge. The polarization surface field propagates in the form of a transient wave at the speed of light and has a sharp front at \(z = ct\). This is shown in Fig. 1. The integrable divergence at the wave front in Eq. (5) is due to the singularity of the linear point-sized charge.

As it should be in the nonrelativistic limit \(v \ll c\), polarization field (5) is like a quasistationary dipole field of two charges (the real charge and its mirror image) at the target surface under the trail of the moving charge with two singular spikes from radiation field running in opposite directions with negligible energy. In the ultrarelativistic limit, this picture changes dramatically. For \(v \gg c\), the radiation field \(E_{y0}^v \ll E_r \sim \lambda/\sqrt{c^2t^2 - z^2}\) dominates the dipole field (compare the red and blue lines in Fig. 1) and spreads away over the surface under the trail of the moving charge. We note that the full surface charge \(\lambda_s = (\int E_{x0}^v dz)/4\pi\) exactly equals the escaping charge \(\lambda\), i.e., \(\lambda_s = -\lambda\). Bearing in mind the laser generation of electrons, we below consider the ultrarelativistic case \(\gamma \gg 1\).

The case of a point-sized source well highlights the difference in target surface charging by slow and fast escaping electrons, which is directly illustrated by the evolution of the normal component of the electric field (Fig. 1). Because the electric field \(E_{x0}^v\) is exactly proportional to the surface charge, it can be clearly seen that the relativistic outgoing electron charge over time forms a positive surface charging in the form of single waves moving apart, in contrast to the charge spot spreading smoothly in the nonrelativistic case. We now show that the same picture basically holds for a finite-sized electron charge generated by a laser in both the normal (pulse width) and transverse (spot size) directions. Assuming that \(|\epsilon|\gamma \gg 1\) in Eq. (2) and using the relations \(B_x = (\omega E_x + 4\pi i j_x)/(k_e)\) and \(E_z = ic/(\omega e)\partial B/\partial x\), we

\[FIG. 1. Vacuum-side surface electric field \(E_{x0}^v(t, z)\) for \(v = 0.1c\) (blue line) and \(v = 0.99c\) (\(\gamma \simeq 7\)) (red line).\]
obtain the approximation

\[
E^v_x = -\frac{4\pi en e^{-k_x x}}{ck_v}, \quad E^\rho_x = \frac{4\pi e n}{\epsilon} \left[ e^{i\omega t - \frac{\epsilon k_x^2 x^2}{2 c^2}} e^{i\omega t - \frac{\epsilon k_x^2 x^2}{2 c^2}} \right],
\]

\[
B^v_y = -\frac{4\pi en}{ck} \left( \frac{\omega}{ck_v} e^{-k_x x - i\epsilon k_x^2 x} \right),
\]

\[
B^\rho_y = -\frac{4\pi en}{ck} \left( \frac{\omega}{ck_v} e^{i\omega t - \frac{\epsilon k_x^2 x^2}{2 c^2}} \right),
\]

\[
E^v_z = \frac{4\pi i en}{ck^2 k_p} \left( 1 + \frac{\omega}{\omega_p} \left( \frac{\omega}{ck_v} - i \right) \right) e^{i\omega t - \frac{\epsilon k_x^2 x^2}{2 c^2}},
\]

\[
E^\rho_z = \frac{4\pi i \omega n}{ck^2 k_p} \left( \frac{\omega}{ck_v} - i \right) e^{i\omega t - \frac{\epsilon k_x^2 x^2}{2 c^2}}.
\]

(6)

To illustrate the generation and evolution of a polarization wave by a finite-sized laser-accelerated electron bunch, we choose the plasma permittivity \(\epsilon = 1 - \omega_p^2/\omega^2 \simeq -\omega_p^2/\omega^2\) and the current source \(en^\rho = \lambda\theta(t)(\theta(t - \tau) - \theta(t - \tau - x)) \exp(-z^2/L^2)/(\lambda t^2/\sqrt{\pi})\). Here, the bunch with \(v \simeq c\) enters the vacuum at \(t = 0\) and leaves the target at \(t = \tau\), and the Heaviside step functions \(\theta(t)\) and \(\theta(t - \tau - x)\) model the laser-heated spot size \(L\) and pulse duration \(\tau\). Using \(en = i\lambda e^{-k^2 L^2/4(1 - e^{i\omega t})}/(\omega^2)\) in Eqs. (6), we studied the evolution of the polarization wave, which is shown in Fig. 2. The amplitude of the induced electric field at the surface initially \((t < \tau)\) increases inside the interaction spot to the maximum value defined by the ratio \(L/c\tau\) and at \(L = c\tau\) is \(\sim 6\lambda/L\). After a pulse terminates \((t > \tau)\), the field separates into two wave bunches propagating at the speed of light in opposite directions along the target surface away from the interaction spot. During propagation, the field amplitude decreases with time as \(1/t\). In accordance with Fig. 2, an electron charge of \(\sim 100 e\) escaping from a spot with \(L \sim 5\mu m\) can generate a surface electric field up to TV/m.

The tangential field \(E^\rho_x \propto 1/\sqrt{|\epsilon|}\) corresponds to the strong plasma current \(j^\rho_x \simeq -i\omega E^\rho_x/(4\pi) \propto \sqrt{|\epsilon|}\) associated with the polarization wave, which behaves similarly to the magnetic field \(B^\rho_y\) in Eq. (6). For \(t > \tau\), we can write the electron current as

\[
j^\rho_x = \frac{i\lambda \omega_{\rho}}{2c\tau} \int \frac{dk}{\pi k} \left[ J_0(ck(t - \tau)) - J_0(ck t) \right] e^{-\frac{k^2 l^2}{L^2}} e^{ikz + \frac{\pi}{4}}.
\]

(7)

This current runs at the speed of light inside the skin layer in the form of two unipolar pulses propagating away from the interaction spot as shown in Fig. [3]. The full linear charge inside these two pulses is exactly equal to the linear charge \(\lambda\) of escaping electrons. Because the polarization-triggered current is excited in a neutral plasma, we interpret it as a charging current unlike the previously assumed charge-neutralizing current \(B\). The polarization wave field is a high-frequency field, whose characteristic frequency \(\omega_s\) can be estimated as \(\omega_s \simeq \min\{\tau^{-1}, c/L\}\).

Because the polarization wave amplitude decreases monotonically \(\propto 1/t\), it eventually drops below the surface wave amplitude, which was initially smaller by the factor \(1/\sqrt{|\epsilon|}\). These amplitudes become comparable at the instant \(t_s \sim \sqrt{|\epsilon|}/\omega_s\), where \(|\epsilon|\) should be evaluated at the frequency \(\sim \omega_s\), or, equivalently, at the distance \(l_s \sim c t_s\). For a plasma target, \(t_s \sim \omega_p^2/\omega_s^2\), and for the metal target, \(t_s \sim \omega_s^{-3/2}\sqrt{\sigma}\), where \(\sigma\) is the electrical conductivity at \(\omega \sim \omega_s\). The distance \(l_s\) increases rather quickly with the laser pulse duration \(\tau\) if \(L/c < \tau\), and for a 1 ps pulse reaches a few cm to several tens of cm. At a large distance \(z > l_s\), the transient polarization wave disappears, and only the surface wave remains. For a plasma target, the latter is described by the analytic expression

\[
E^\rho_x = \frac{4\sqrt{\pi} \lambda}{\omega_p \tau L} \left( e^{-\frac{(ct_\pm z)^2}{L^2}} - e^{-\frac{(ct - \tau \pm z)^2}{L^2}} \right) \theta(\pm z).
\]

(8)

We note that in deriving Eq. (5), we neglect the imaginary part of \(\epsilon\). A more accurate calculation shows that the surface wave propagates with a velocity slightly different from the speed of light in the form of a bipolar
pulse (a two-pi pulse) unlike the polarization wave as seen in Fig. 2.

Our theory ignores the possible permanent loss of a small fraction of the laser-accelerated electrons from a target hot spot. The theoretical description of the related charging surface field and current still remains unclear, although such a postulation has been attributed to the experiment \(^3\). To clarify whether a charge-neutralizing current could be detectable, we performed three-dimensional FDTD numerical simulations using the code VSim. To advance the study of a transient field propagating along the target at the speed of light, we used an escaping bunch with a small velocity \(v_0\). In accordance with the analytic theory, this simulation result in Fig. 4, which agree well with the theoretical model for \(\lambda_0 = 3\) \(\mu\)m shown by the PIC result. This field decreases similarly to \(1/z\) and predictably drops to the multi-GV/m level at distance of \(\sim1\) cm. In this regard, we note that electromagnetic pulses propagating at the speed of light with a strong electric field (of the order of GV/m) have already been measured at a cm distance from the hot spot \(^5\). As a final note, we emphasize that a small fraction of the laser-accelerated hot electron population expands along the target surface from the vacuum side because it is well held by the surface fields (see \(^2\) \(^3\) \(^5\)).

In summary, our theory sheds light on the physical

\[
e_\varepsilon \left( x, y, z \right) = \frac{1}{v_0^2 - c^2} \left( 1 + 4\pi i \sigma \left( \omega \right) / \nu \right) \frac{\sigma_0}{\nu}, \quad \sigma = \sigma_0 / \left( 1 - i \omega / \nu \right), \quad \sigma_0 = 10^{18} \text{s}^{-1} \quad \nu = 10^{13} \text{ s}^{-1}.
\]

The target occupied a half-space \(x < 75 \mu\)m in the full simulation box \(0 < x < 300 \mu\)m, \(-300 \mu\)m < \(y < 300 \mu\)m, and \(-300 \mu\)m < \(z < 300 \mu\)m. The grid cell size was \(1 \mu\)m and the time step was \(1 \text{ fs}\).

In contrast to the analytic model, we assumed full charge neutrality in the FDTD simulation resulting in target charging during electron escape. We show our simulation results in Fig. 4 which agree well with the analytic model showing that the polarization wave is the main contributor to the surface electromagnetic field.

In accordance with the analytic theory, this simulation shows that a change in the type of dielectric permittivity has only a minor effect on the surface field distributions until \(|\epsilon| \gg 1\). In a similar simulation where we used an escaping bunch with a small velocity \(0.2\), we observed a barely smooth spreading of the field rather than a wavelike structure, which also corresponds to the analytic theory.

\[
e_\varepsilon \left( x, y, z \right) = \frac{1}{v_0^2 - c^2} \left( 1 + 4\pi i \sigma \left( \omega \right) / \nu \right) \frac{\sigma_0}{\nu}, \quad \sigma = \sigma_0 / \left( 1 - i \omega / \nu \right), \quad \sigma_0 = 10^{18} \text{s}^{-1} \quad \nu = 10^{13} \text{ s}^{-1}.
\]

The maximum surface electric field initially reaches \(\sim15 \text{ TV/m}\) in agreement with our theoretical model for \(\lambda \sim 2 \mu\)C/\(\mu\)m shown by the PIC result. This field decreases similarly to \(1/z\) and predictably drops to the multi-GV/m level at distance of \(\sim1\) cm. In this regard, we note that electromagnetic pulses propagating at the speed of light with a strong electric field (of the order of GV/m) have already been measured at a cm distance from the hot spot \(^5\). As a final note, we emphasize that a small fraction of the laser-accelerated hot electron population expands along the target surface from the vacuum side because it is well held by the surface fields (see \(^2\) \(^3\) \(^5\)).

In summary, our theory sheds light on the physical
mechanism of the generation and propagation of a transient electromagnetic pulse and a lateral current in the wave form along the target surface at the speed of light. The proposed mechanism is associated with fast electric polarization of a high-conductivity target during ejection of a laser-driven electron bunch from a target into a vacuum. Supported by the results of the developed theory and two simulation models, the mechanism proposed and studied here might be important for a deeper understanding of the experiments, which have already shown the extremely strong ultrafast charging of a solid irradiated by a high-intensity laser and the propagation of the corresponding transient laser-triggered current along the target surface [2–7].

This research was supported by the Russian Science Foundation.

[1] T. Nakamura, S. Kato, H. Nagatomo and K. Miina, Phys. Rev. Lett. 93, 265002 (2004).
[2] P. McKenna, D. C. Carroll, R. J. Clarke, R. G. Evans, K. W. D. Ledingham, F. Lindau, O. Lundh, T. McCanny, D. Neely, A. P. L. Robinson, L. Robson, P. T. Simpson, C.-G. Wahlstro, M. Zepf, Phys. Rev. Lett. 98, 145001 (2007).
[3] K. Quinn, P. A. Wilson, C. A. Cecchetti, B. Ramakrishna, L. Romagnani, G. Sarri, L. Lancia, J. Fuchs, A. Pipahi, T. Toncian, O. Willi, R. J. Clarke, D. Neely, M. Notley, P. Gallegos, D. C. Carroll, M. N. Quinn, X. H. Yuan, P. McKenna, T. V. Liseykina, A. Macchi, and M. Borghesi, Phys. Rev. Lett. 102, 194801 (2009).
[4] G. Sarri, A. Macchi, C. A. Cecchetti, S. Kar, T. V. Liseykina, X. H. Yang, M. E. Dieckmann, J. Fuchs, M. Galimberti, L. A. Gizzi, R. Jung, I. Kourakis, J. Osterholz, F. Pegoraro, A. P. L. Robinson, L. Romagnani, O. Willi, and M. Borghesi, Phys. Rev. Lett. 109, 205002 (2012).
[5] S. Kar, H. Ahmed, R. Prasad, M. Cerczez, S. Brauckmann, B. Auran, G. Cantono, P. Hadjisolomou, C. L. S. Lewis, A. Macchi, G. Nersisyan, A. P. L. Robinson, A. M. Schroer, M. Swantusch, M. Zepf, O. Willi and M. Borghesi, Nat. Commun. 7, 10792 (2016).
[6] R. Pompili, M. P. Anania, F. Bisesto, M. Botton, E. Chiadroni, A. Cianchi, A. Curcio, M. Ferrario, M. Galletti, Z. Henis, M. Petrarca, E. Schleifer, A. Zigler, Sci. Rep. 8, 3243 (2018).
[7] S. Tokita, S. Sakabe, T. Nagashima, M. Hashida, S. Inoue, Sci. Rep. 5, 8268 (2015).
[8] A.S. Kuratov, A.V. Brantov, Yu.M. Aliev, V.Yu. Bychenkov, Quantum Electronics 48, 653, (2018).
[9] H. Hamster, A. Sullivan, S. Gordon, W. White, and R. W. Falcone, Phys. Rev. Lett. 71, 2725 (1993).
[10] A. Gopal, S. Herzer, A. Schmidt, P. Singh, A. Reinhard, W. Ziegler, D. Bronnem, A. Karmakar, P. Gibbon, U. Dillner, T. May, H-G. Meyer, and G. G. Paulus, Phys. Rev. Lett. 111, 074802 (2013).
[11] V. L. Ginzburg and V. N. Tsytovich, Phys. Reports 49, 1 (1979).
[12] V. L. Ginzburg and I. M. Frank, JETP (USSR) 16, 15 (1946).
[13] V. Ya. Eidman, Izvestiya VUZ, Radiofizika 8, 188 (1965).
[14] B.M. Bolotovskii and A.V. Serov, Physics Uspekhi 52, 487 (2009).
[15] H. Nakajima, S. Tokita, S. Inoue, M. Hashida, and S. Sakabe, Phys. Rev. Lett. 110, 155001 (2013).