Optical-to-THz Wave Conversion via Excitation of Plasma Oscillations in the Tunneling-Ionization Process

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A new ("linear-parametric") mechanism of a direct conversion of an ultrashort laser pulse into terahertz radiation is suggested. The conversion is due to the ionization-induced excitation and the subsequent electromagnetic emission of the superluminous polarization wave created by the axicon-focused laser pulse. For a few-cycle pulse with an optimum carrier-envelope phase, the considered mechanism is found to be much more effective than the alternative one based on the excitation of plasma oscillations in the laser wakefield by the ponderomotive force and able to provide THz radiation of the gigawatt power level with the use of moderate optical intensity ($\sim 10^{14}$–10^{15} W/cm^2).

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The use of collective properties of a laser-produced plasma underlies a number of recently proposed methods of THz waves generation. In particular, many projects actively worked out at present are based on the phenomena of ponderomotive-force-induced excitation (and subsequent electromagnetic radiation) of the Langmuir oscillations [1,2]. This excitation mechanism is nonlinear in its physical nature and can provide sufficiently large ($\sim$10 MW) power of THz radiation only at very high intensity of the exciting laser pulses ($\sim 10^{19}$ W/cm^2 at pulse duration $\sim$100 fs) [2].

In this Letter, we would like to draw attention to another laser-plasma method of THz wave generation based on the ionization-induced excitation of the natural plasma oscillations during the gas breakdown produced by an ultrashort (few-cycle) laser pulse. The initial push for the natural plasma oscillations is caused, in this case, directly by the optical electric field imparting a large constant velocity component (depending on the field phase) to the newly born electrons [3]. The ionization process apart, this excitation mechanism by itself is linear (or "linear-parametric") with respect to the pumping (exciting) optical field and, therefore, can ensure much more efficient conversion of this field into the radiating plasma oscillations. The main obstacle for its realization using a multicycle (narrow-spectrum) laser pulses is a very strong spatiotemporal spread in the initial phase of electron oscillations resulting in a drastic decrease in the ordered electron current (responsible for the coherent THz radiation). This obstacle, however, becomes not very significant for the few-cycle ($\sim$1–10 fs) pulses with a controlled phase of the electric field. As we will see, the ionization-induced plasma oscillations in the wakefield of such a pulse focused by the axicon lens can provide a very high (GW) level of THz radiation at a much lower optical intensity ($\sim 10^{14}$–10^{15} W/cm^2) than the ponderomotive-force-induced ones. Thus, recent successes in the generation and phase control techniques of the few-cycle pulses [4,5] offer new possibilities in THz generation methods and put the study of plasma oscillations excited in the ionization process by the few-cycle laser pulses among urgent problems of laser-plasma physics.

The excitation mechanism under consideration is related to that studied theoretically and experimentally in Refs. [6–8], where the natural plasma oscillations were excited by extraneous (static or microwave) fields, and the fs laser pulse served for plasma creation only. In this case, the maximum permissible strength of the external (pumping) field is determined by the threshold of gas breakdown, which limits strongly the amplitude of oscillations and the accessible value of THz radiation. The method of generation we suggest here is free from this limitation as both the gas ionization and excitation of plasma oscillations are produced by the same strong laser electric field (comparable with the atomic one by the order of magnitude).

A rather promising generation scheme allowing creation of a sufficiently long radiating plasma column is based on the use of axicon focusing [7–9]. In the framework of this generation scheme, we analyze here the radiation originating from a fast transverse polarization wave that is excited behind the front of a superluminous breakdown wave (see Fig. 1) in a narrow near-axial region of the conical lens (axicon). The linear polarized laser pulse focused by this lens into a so-called Bessel wave beam propagates without divergence along the symmetry axis z at the distance $L = b/\tan\vartheta_0$ determined by the wave beam radius b at the entrance of the lens and the focusing angle ϑ_0. The velocity of the ionization front created by the laser pulse in the segment $\Delta z = L$ is equal to the phase velocity of the light wave, $V_i = c/\cos\vartheta_0 > c$. The transverse ($x$-directed) electric field $\mathbf{E}_L$ of the laser pulse (as well as the density N of the field-created plasma) can be approximately presented as functions of the delaying time $t' = t - (z/V_i)$ and the distance r from the axis z:

$$\mathbf{E}_L = \mathbf{e}_x E_0 \cos(\omega_L t' + \psi)\exp\left(-\frac{t'^2}{\tau^2}\right) J_0(k_\perp r), \qquad (1)$$

where E_0 is the envelope maximum, ω_L is the laser carrier

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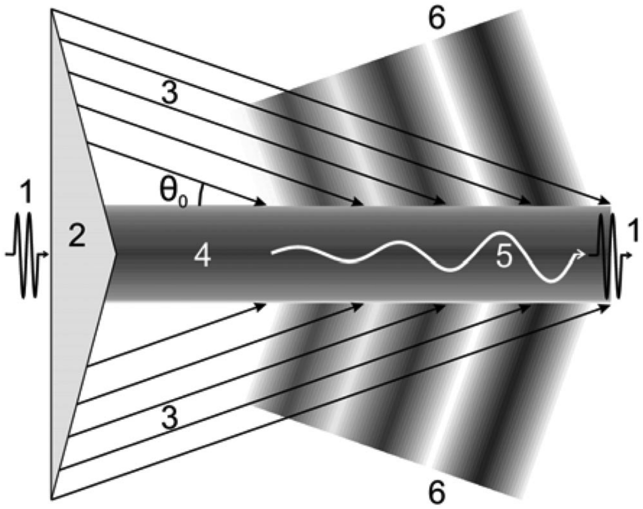

FIG. 1. Schematic of the optical-to-THz wave conversion; 1—laser pulse; 2—axicon lens; 3—refracted rays; 4—pulse-created plasma; 5—laser-field-induced transverse polarization wave; 6—THz radiation fronts.

frequency, ψ is the carrier-envelope (CE) phase, $\tau = \tau_p/\sqrt{2\ln 2}$, τ_p is the pulse duration (full-width at half-maximum), $J_0(\xi)$ is the zero-order Bessel function, $k_\perp = (\omega_L/c)\sin\vartheta_0$. The focusing angle is assumed to be small ($\vartheta_0 \ll 1$), so that the longitudinal component of the electric field can be neglected. Strictly speaking, Eq. (1) is correct when $\omega_L\tau \gg 1$ only, but it can also be used as an approximate one in the case of interest, $\omega_L\tau \sim 1$.

Time variations of the plasma density and free-electron current density $\mathbf{j}$ in the case of single-electron tunneling ionization are governed by the equations [3,10]

$$\frac{\partial N}{\partial t'} = (N_g - N)f(|\mathbf{E}|), \qquad \frac{\partial \mathbf{j}}{\partial t'} + \nu\mathbf{j} = \frac{\omega_p^2}{4\pi}\mathbf{E}, \tag{2}$$

where $\mathbf{E} = \mathbf{E}_L + \mathbf{E}_p$ and $\mathbf{E}_p$ is the field created by the charge and current of plasma, the function f is defined by the kind of the gas, N_g is the undisturbed gas density, $\omega_p = [4\pi e^2 N/m]^{1/2}$ is the plasma frequency, ν is the effective collision frequency of electrons with heavy particles.

The radius a of the pulse-produced inhomogeneous plasma cylinder, the maximum plasma frequency $\omega_{pm} = \omega_p$ at $N = N_g$, and the pulse duration τ are assumed to satisfy the conditions

$$\omega_{pm}a \ll c, \qquad \omega_{pm}\tau \ll 1, \qquad \omega_{pm}L \gg c, \tag{3}$$

which make it possible to find the field $\mathbf{E}_p$ and electron current of the natural oscillations (with the frequency $\sim\omega_{pm}$) in each given section $z = \text{const}$ by solving the two-dimensional quasielectrostatic initial problem disregarding the longitudinal inhomogeneity of the resultant plasma. In this approach, the z coordinate appears in the solution only as a parameter determining approximately the time instant $t_i = z/V_i$ of plasma creation and electron current initiation.

Under conditions (3), all the evolution processes discussed above can be divided into two stages. At the short first stage, during the pulse passing the given point z (within the time interval $|t'| \sim \tau \ll \omega_{pm}^{-1}$), the plasma density and electron current density are the fast-varying time functions. Since the radius a is determined by the transverse scale of the Bessel beam ($a \sim k_\perp^{-1}$), the first condition (3) is equivalent to $\omega_{pm} \ll \omega_L \sin\vartheta_0$, which is the condition of negligible influence of the plasma on the optical field [8]. Thus, at the first stage, the evolution Eqs. (2) can be integrated with $\mathbf{E} = \mathbf{E}_L$:

$$N(r, t') = N_g\left[1 - \exp\left(-\int_{-\infty}^{t'} f(|\mathbf{E}_L(r, \tau')|)d\tau'\right)\right], \tag{4}$$

$$\mathbf{j}(r, t') = \frac{e^2}{m}\int_{-\infty}^{t'} N(r, \tau')\mathbf{E}_L(r, \tau')d\tau'. \tag{5}$$

In the last equation, we have neglected the collision losses assuming that $\nu \ll \omega_L$.

At the second stage, after laser pulse passing, the natural transverse oscillations take place against the background of the steady-state density profile $N_s(r) = N(r, \infty)$ determined formally by solution (4) at $t' = \infty$. In the case we consider, as it follows from the results of numerical calculation (see Ref. [8] and some results presented below), this profile combines a rather wide plateau ($N = N_g = \text{const}$) with a comparatively narrow transition region where the density reduces smoothly from N_g to zero. For such profiles, the natural oscillations of the transverse dipole moment P (per unit of length along the z axis) are described approximately by an oscillatory equation [8]:

$$\frac{d^2P}{dt'^2} + 2\gamma_i\frac{dP}{dt'} + \omega_c^2 P = 0. \tag{6}$$

Here, $\omega_c = \omega_{pm}/\sqrt{2}$ is the frequency of the so-called geometrical resonance of the plasma cylinder; the internal damping constant $\gamma_i = \nu/2 + \gamma_l$ is determined by the electron collisions and resonant losses in the transition layer; on the assumption that $\gamma_i \ll \omega_c$, we have (see also Refs. [8,11]) $\gamma_l = \omega_c l/a$, $l = |N/\nabla N|_{N=N_c}$, where $N_c = m\omega_c^2/(4\pi e^2) = N_g/2$.

The initial conditions for Eq. (6) are determined by the current distribution at the end of the first stage, that is by Eq. (5) at $t' = \infty$. In view of the inequality $\omega_{pm}\tau \ll 1$ we can take, as the initial point of the second stage, the time instant $t' = 0$ assuming

$$P(0) = 0, \qquad \frac{dP}{dt'}(0) = \int_0^\infty j_x(r, \infty)2\pi r dr \tag{7}$$

and presenting the solution of Eq. (6) in the form

$$P = P_0 \sin\omega_c t' \exp(-\gamma_i t')\theta(t'), \tag{8}$$

where $P_0 = \omega_c^{-1} dP/dt'(0)$, $\theta(t' < 0) = 0$, and $\theta(t' > 0) = 1$. The efficiency of the process in hand can be characterized by the dimensionless excitation factor $K = k_\perp^2 P_0 \omega_L/(E_0 \omega_c)$ written in the dimensionless variables $t' \to \omega_L t'$, $r \to k_\perp r$, $N \to N/N_g$, $E_L \to E_L/E_0$ as

$$K = \int_{-\infty}^{\infty} dt' \int_0^\infty rNE_L dr. \tag{9}$$

This factor, depending on the duration, intensity, and phase structure of the pulse, was calculated numerically based on Eqs. (1) and (4) with the function f corresponding to the tunneling ionization of the hydrogen atoms: $f = 4\Omega(E_a/|\mathbf{E}|)\exp(-2E_a/3|\mathbf{E}|)$, where $\Omega = 4.13 \times 10^{16}\ \mathrm{s}^{-1}$ and $E_a = 5.14 \times 10^9$ V/cm are the atomic units of the frequency and field strength. The results of the calculation are presented for $\omega_L = 2.4 \times 10^{15}\ \mathrm{s}^{-1}$ (the central wavelength $\lambda_L \approx 800$ nm) in Figs. 2–4. Given the pulse duration τ_p and envelope maximum E_0, the excitation factor K (see Fig. 2) proves to be a sign-alternating function of the CE phase ψ [$K(\psi) = -K(\psi + \pi)$] achieving a maximum absolute value $K_m = |K|_{\max}$ at some point $\psi = \psi_m$. The curves $K_m(\tau_p)$ are shown at different $E_0 \ll E_a$ in Fig. 3. The values K_m have maxima at $\tau_p \approx 0.9$ fs ($\omega_L \tau_p \approx 2$) and tend to zero with both $\tau_p \to 0$, and $\tau_p \to \infty$. Time dependencies of the laser field $E(0, t')$ and plasma density $N(0, t')$ at $\psi = \psi_m$ and the density profiles $N(r, t')$ are illustrated for two-cycle pulse ($\omega_L \tau_p = 12$) in Fig. 4. In this example, the steady-state profile parameters are $k_\perp a \approx 1.5$, $k_\perp l \approx 0.25$, and the collisionless damping constant is $\gamma_l/\omega_c = l/a \approx 0.16$.

We note that the approach we use here based on Eqs. (6)–(9) is more adequate for the calculation of laser-pulse-induced polarization and optimum phase ψ_m than the simple model used before in Ref. [12]. The reason is that the resultant dipole moment of plasma is determined by the transverse asymmetry of the initial (arising immediately after passing of the pulse) electron current, and the value that determines this asymmetry (at least, in the realistic case of interest $\nu \ll \omega_L$) is precisely the excitation factor (9) but not the difference of the electron densities arising at the positive and negative field phase as it was supposed in [12].

The traveling fast polarization wave described by Eq. (8) emits in the directions forming the angle ϑ_0 with z axis. The frequency ω_c of the emitted radiation is determined by the neutral gas density N_g; in particular, $f_c = \omega_c/2\pi = 3$ THz at $N_g \approx 2.2 \times 10^{17}\ \mathrm{cm}^{-3}$ (the gas pressure $p \approx 6$ Torr). As is easily shown by calculating the electromagnetic field in the far zone of the given linear radiator of the length $L \gg c/\omega_c$ and generalizing Eq. (6) with inclusion of radiative losses, the angular width of the emission maximum $\Delta\vartheta$, total peak power Π and energy W of radiation for the case of weak damping are given by the expressions

$$\Delta\vartheta = \frac{c}{\omega_c L \sin\vartheta_0}, \qquad \Pi = 2\gamma W, \qquad W = W_0 \frac{\gamma_r}{\gamma}. \tag{10}$$

Here, $\gamma_r = \pi\omega_c^3 a^2/8c^2$ is the radiative damping constant, $\gamma = \gamma_i + \gamma_r$ is the total line width including the internal and radiative losses, $W_0 = P_0^2 L/a^2$ is the emitted energy achievable at $\gamma_i \ll \gamma_r$. Given the parameters K, $k_\perp a$, γ_r/γ, and $\vartheta_0 \ll 1$, the energy conversion efficiency η is conveniently written as

$$\eta \equiv \frac{W}{W_L} = \frac{16\pi\gamma_r}{(k_\perp a)^2 \gamma} \frac{K^2}{\omega_L \tau_p} \frac{\omega_c^2}{\omega_L^2 \vartheta_0^2}, \tag{11}$$

where $W_L = (c^2/16\pi\omega_L)\tau_p L E_0^2$ is the energy of the laser pulse.

The radiation proves to be maximal at the boundary of the region of the quasistatic description ($a \ll c/\omega_c$) where the obtained results enable an approximate estimates [8]. Under the optimum conditions, taking $\omega_c a/c \sim 1$, $\gamma \sim \gamma_r$, $k_\perp a \approx 1.5$, and $K^2 \approx K_m^2$ we find

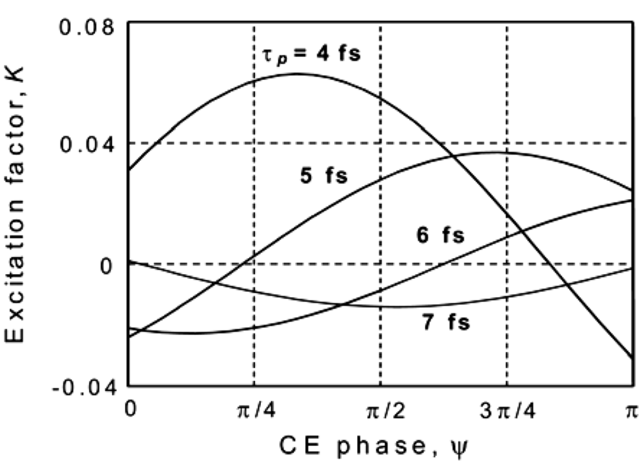

FIG. 2. Excitation factor K as a function of the CE phase ψ for the pulse envelope maximum $E_0/E_a = 0.15$, central wavelength 800 nm, and different pulse duration τ_p.

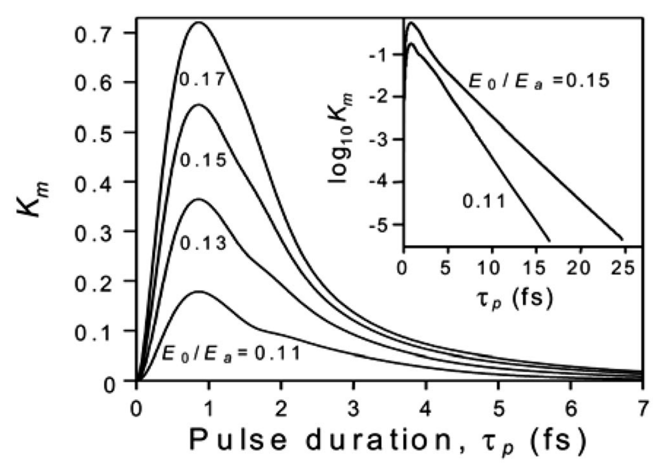

FIG. 3. Maximum excitation factor $K_m = |K|_{\max} = |K(\psi_m)|$ as a function of the pulse duration τ_p for different E_0/E_a.

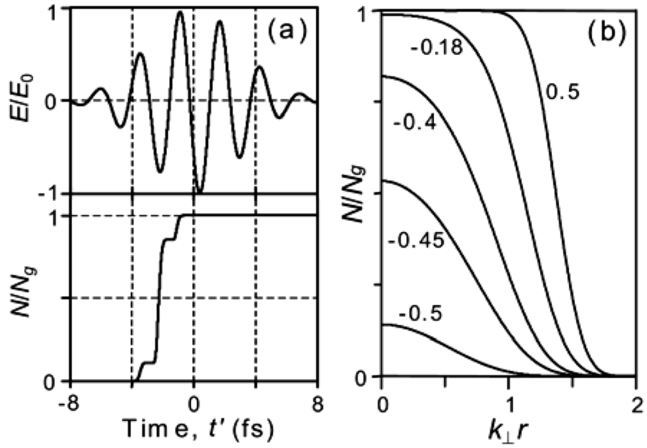

FIG. 4. Evolution of the optical electric field and plasma density for the optimum CE phase $\psi = \psi_m = 0.7\pi$, $E_0/E_a = 0.15$, and pulse duration $\tau_p = 5$ fs ($\omega_L \tau_p = 12$); (a) the functions $E(t')/E_0$ and $N(t')/N_g$ at the beam axis ($r = 0$); (b) density profiles $N(r)/N_g$ at different time instants (numerals by the curves show the values t'/τ_p); the second slight maximum ($N/N_g < 0.025$) formed near the second maximum of the Bessel function ($k_\perp r \approx 3.8$) is not shown here.

$$\Pi \sim \omega_c W, \qquad \eta \sim \frac{10K_m^2}{\omega_L \tau_p}, \qquad \vartheta_0 \sim \frac{1.5\omega_c}{\omega_L}. \tag{12}$$

In particular, in the example considered above, at the parameter values $\omega_L \tau_p = 12$, and $E_0/E_a = 0.15$, the 3 THz radiation can be generated with the efficiency $\eta \sim 10^{-3}$, energy $W \sim 1$ μJ, and power $\Pi \sim 10$ MW. The above pulse parameters correspond to the few-cycle laser pulses realized in the experiment [4] ($\tau_p = 5$ fs, $W_L = 0.5$ mJ), when focused to the Bessel beam with the axial intensity $I = 7.9 \times 10^{14}$ W/cm^2 and caustic dimensions $L = 20$ cm, $a \sim 10$ μm ($\vartheta_0 \sim 10^{-2}$, $b \sim 0.2$ cm). The emitted THz energy in this example are of the same order as the one achieved in Ref. [2] for the ponderomotive-force-induced mechanism of the electron acceleration at the laser intensity $I \sim 10^{19}$ W/cm^2 and pulse energy $W_L \sim 1$ J, whereas the conversion efficiency in our case turns out to be several orders higher. The significant gain in the emitted energy and efficiency could be achieved at the further success in generation of the high intensity few-cycle pulses: the pulse with $W_L = 0.5$ mJ and $\tau_p = 1$ fs could be converted into a 3 THz wave of 10 GW power with the efficiency $\eta \sim 1$ (in the framework of the above-used estimations). These estimations could be improved by taking into account the laser pulse energy losses in plasma and changes of the pulse form due to nonmonochromaticity and dispersion in the axicon material, as well as THz-field corrections for the finite $\omega_c a/c$ parameter value.

In summary, using a very short ("few-cycle") laser pulse of comparatively modest intensity projected through an axicon lens into a rarefied gas, the authors show that one can generate a super-light-speed coherent pulse current in the plasma column created by the laser pulse. The coherent transverse traveling-wave current is formed by the ionization electrons with the coherent flow that is the result of the ionization via the laser field, in the direction of the incident laser polarization. The result is an emission in what might be called a relativistic Mach cone, at the plasma column's dipole transverse resonance frequency. This is typically THz, being about the column's plasma frequency (proportional to the square root of the gas density). With the laser pulse length not much longer than the optical field period, the conversion efficiency at the modest intensity suitable for ionization is several orders of magnitude higher than the ponderomotively-driven mechanisms usually considered (these last usually require very intense pulses for an acceptable level of the THz emission). Once the available laser pulses are sufficiently short, the prospects for very interesting experiments are bright.

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