Bulk-to-surface-wave self-conversion in optically induced ionization processes

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Nonlinear time evolution of a *p*-polarized wave mode with inhomogeneous transverse structure producing tunnel ionization of a gas is investigated by numerical simulation and theoretical analysis. A phenomenon of trapping of electromagnetic radiation via its adiabatic conversion into surface waves guided by the field-created plasma structure is found out numerically. This process is accompanied by significant frequency downshifting of the electromagnetic radiation. The underlying physical mechanism is explained using a simple theoretical model. The described phenomena may play significant role in the self-channeling and frequency tuning of intense ($\sim 10^{14}-10^{18}$ W/cm^2) laser pulses in dense gases.

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I. INTRODUCTION

In the last decade, the advances in generation of ultrashort (10–100 fs) and highly intense ($>10^{14}$W/cm^2) laser pulses have drawn attention to many new basic aspects in the physics of field-matter interaction and stimulated the elaboration of important applications of laser-produced plasmas. Among these applications are x-ray lasers, laser-driven particle acceleration, plasma-based sources of tunable radiation, and creation of elongated plasma waveguiding structures (see, e.g., [1–4] and references therein). One of the key physical factors, which determines the dynamics of intense short laser pulses in a medium, is the ionization nonlinearity of the medium [5–10]. In the case of strong fields, ionization-induced nonlinear mechanisms are practically inertialess and manifest themselves already at the initial stage of the plasma creation process. Thereby, they determine in many cases the features of the combined field-plasma structure produced at the further stages of the process. Two of the most interesting and significant manifestations of the ionization nonlinearity are the phenomena of frequency self-shifting and self-channeling of the laser radiation.

The phenomenon of frequency shifting is well known for the electromagnetic wave in media with specified time variations of parameters [11]. The effect of frequency self-upshifting (or blue-shifting) of an intense electromagnetic pulse propagating through a gas and producing gas ionization was first demonstrated by Yablonovitch [12] and observed later in several experiments both in optical [13] and microwave region [14]. This effect has potential applications in tuning laser radiation and as a diagnostic tool for structures moving with relativistic velocities in plasmas [15]. Theoretical description of the effect for the case of a small perturbation of the pulse spectrum was treated in terms of weak self-phase modulation of the pulse [11]. Some peculiarities of the effect for focused electromagnetic pulses were considered within the framework of the paraxial beam (parabolic) approximation [16]. The strong adiabatic frequency self-upshifting of one dimensional (1D) and focused (2D) pulses producing gas ionization was considered both within the WKB approach (modified by including the fast oscillations of the tunnel ionization rate) and by numerical integration of the wave equation [17–20]. Essentially, a plane electromagnetic wave or paraxial wave beam propagating in a plasma with time-growing density and approximately constant electron collision frequency can only be frequency upshifted. The frequency downshifting of a plane wave is possible only if plasma decay processes prevail over the ionization or due to the fast rise of the collision frequency [21].

The self-channeling of ultrashort electromagnetic pulses is of particular interest because some important applications of superstrong fields require formation of rather long laser-produced plasma structures for interaction with optical radiation over distances of many Rayleigh lengths. Several mechanisms of self-channeling were discussed. At high intensities ($\geqslant 10^{18}$ W/cm^2 at near-infrared wavelengths) self-channeling can occur as a result of two effects: the relativistic modification of electron mass in the laser field and the reduction of the electron density on axis due to the expulsion of electrons by laser pondermotive force [22–25]. Experimentally observed extra-long waveguides produced by a few millijoules, 100-fs laser pulses at ionization of atmospheric air [26,27] have been interpreted as plasma structures having a core where the defocusing ionization nonlinearity prevails and an outer cladding where the focusing Kerr nonlinearity generates positive-in-sign variations of the refractive index and hence keeps the radiation from divergence [26–29]. More surprising regimes of self-channeling of ionizing radiation without any focusing nonlinearity have been found out as well. Recently, it was demonstrated in Ref. [10] that short laser pulse self-guiding over distances of many Rayleigh lengths can be achieved as a result of the trapping of a leaky mode in a plasma channel produced by field-induced ionization in the saturation regime. Another mechanism of self-channeling of ionizing radiation in an unbounded gaseous medium was predicted as early as in 1983 [30]. It was demonstrated analytically that stationary self-sustained surface-waveguiding plasma structure can exist. In this structure a surface wave produces plasma slab that then guides the surface wave. Such structure is a rare example when the field and isotropic plasma have a localizing effect on each other and, as a result, they are concentrated in the same spatial region. This situation qualitatively differs, for example, from the case of self-channeling due to focusing (pondermotive-induced) nonlinearity where concentration of the field on axis is accompanied by the reduction of electron density on the axis. However, results of Ref. [30] are restricted by the framework of a rather special ionization model that is not quite appropriate for the case of laser induced plasma production. Moreover, these results do not describe the dynamics of the structure formation process. Meanwhile, the investigation of the processes resulting in the formation of such type self-sustained structures are of great interest both from the standpoint of basic theory of nonlinear waves and applications to the problems of laser plasma production and transmittance of electromagnetic energy through an ionized medium.

In the present paper, the nonlinear time evolution of two intersecting high-intensity plane electromagnetic waves producing ionization of a gas is investigated by numerical simulation and theoretical analysis. Tunnel mechanism of ionization is considered. The process of formation of self-sustained plasma waveguides accompanied by conversion of the initial bulk electromagnetic fields into surface waves localized near the plasma slabs is traced numerically (see also Ref. [31]). It is demonstrated that during this process the electromagnetic radiation suffers significant frequency downshifting. Thus, both above-mentioned nonlinear phenomena (self-channeling and frequency self-shifting) turn out to be tightly interconnected here. The underlying physical mechanism is explained using a simple theoretical model. The described phenomena open new aspects in the physics of self-channeling and frequency self-shifting of intense ($\sim 10^{14} - 10^{18}$ W/cm^2) laser pulses in dense gases.

Recently, the effect of conversion of electromagnetic radiation incident on a time-varying plasma structure—plasma half-space [32] or plasma slab [33]—into frequency downshifted surface modes of the structure was found. The case of instant growth of plasma density in time due to effect of an external ionizing factor was considered. Practically, the approximation of rapid plasma creation means that the rise time of the plasma is much shorter than the period of the incident electromagnetic wave. In this case, transient fields are created after the plasma density shift. The transient fields generate frequency downshifted guided modes that propagate along the waveguiding structure. It should be stressed that in the present paper the opposite case of slow, compared to the wave period, variation of plasma density in time is considered. It means that the phenomenon of conversion, we report here, has adiabatic character and, therefore, differs principally from the effect described in Refs. [32,33]. Moreover, we consider the nonlinear problem when plasma density variation is provided by the electromagnetic wave itself, unlike the linear case of Refs. [32,33].

The paper is organized as follows. In Sec. II we present the basic equations and discuss the adiabatic approach. Results of numerical simulation are given in Sec. III. In Sec. IV we explain the underlying physics using a simple theoretical model. Section V gives our conclusion.

II. FORMULATION OF THE PROBLEM—BASIC EQUATIONS

We consider an initial value problem for nonlinear self-consistent evolution of electromagnetic fields and plasma

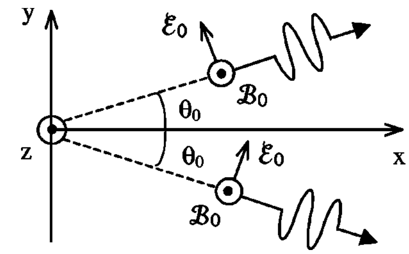

FIG. 1. Geometry of the problem: Initially, two identical plane electromagnetic waves of *p*-polarization intersect in a neutral gas.

density created by the fields via tunnel ionization of a gas. At the initial instant $t=0$ plasma density is set equal to zero and electromagnetic field is a superposition of two intersecting under the angle $2\theta_0$ plane waves that in the absence of ionization would have frequency ω_0 and amplitude of the magnetic field $\mathcal{B}_0$ (see Fig. 1). The magnetic and electric field components if being undisturbed by the arising plasma are given by

$$\mathcal{B}_z = 2\mathcal{B}_0 \cos(g_0 y)\cos(\omega_0 t - h_0 x),$$
$$\mathcal{E}_y = \mathcal{B}_z \cos\theta_0, \tag{1}$$
$$\mathcal{E}_x = -2\mathcal{B}_0 \sin\theta_0 \sin(g_0 y)\sin(\omega_0 t - h_0 x),$$

where $h_0 = (\omega_0/c)\cos\theta_0$, $g_0 = (\omega_0/c)\sin\theta_0$. We use Gaussian system of units where the amplitude of the plane wave electric field is $\mathcal{E}_0 = \mathcal{B}_0$. Due to interference the intersecting plane waves form the wave that is standing along the y axis and running in x direction. Ionization of the gas, beginning at the moment $t=0$, depends on the absolute value of the electric field $\mathcal{E} = \sqrt{\mathcal{E}_x^2 + \mathcal{E}_y^2}$. Therefore, time-varying plasma density $n(x,y,t)$ will be periodic along y axis with period $L = \pi/2g_0$. It allows us to restrict consideration by the spatial interval $0<y<L$.

To describe the process of tunnel ionization of the neutral molecules of the gas we use the well-known static expression for the electron production rate from the lower state of the hydrogen atom [34]

$$\frac{\partial n}{\partial t} = w(\mathcal{E},n) = 4\Omega(N_g - n)\frac{\mathcal{E}_a}{\mathcal{E}}\exp\left(-\frac{\mathcal{E}_a}{\mathcal{E}}\right), \tag{2}$$

where $\Omega = me^4/\hbar^3 = 4.16\times 10^{16}\ \mathrm{s}^{-1}$ and $\mathcal{E}_a = m^2 e^5/\hbar^4 = 5.14\times 10^9$ V/cm are the characteristic atomic frequency and atomic field strength, respectively, N_g is the initial (i.e., before the beginning of ionization processes) neutral molecule density. Equation (2) is applicable to describe ionization in a laser field of frequency ω_0 and amplitude $\mathcal{E}_0$ when the following conditions are satisfied:

$$\omega_0 \ll \Omega, \quad I \ll W_q, \quad \mathcal{E}_0 \ll \mathcal{E}_a, \tag{3}$$

where $W_q = e^2\mathcal{E}_0^2/2m\omega^2$ is the quiver energy of electrons and I is the energy of atom ionization. The expression (2) for $\partial n/\partial t$ gives qualitatively correct results for excited complex atoms. Indeed, this expression incorporates the general characteristic features of the tunnel ionization—fast growth of

$\partial n/\partial t$ with $\mathcal{E}/\mathcal{E}_a$ at small $\mathcal{E}/\mathcal{E}_a$ and saturation at $\mathcal{E}/\mathcal{E}_a \sim 1$. More accurate analysis may be carried out by using the formulas for ionization rates from Ref. [35]. The tunnel mechanism prevails over the impact one ($w \gg \nu_i n$; ν_i is the ionizing collision frequency) in the parameter range $\ln[\Omega N_g/(\nu_i n)] > 1 + (\Omega/\omega_0)(I/W_q)$.

At sufficiently low ionization rate the stated problem may be solved within the adiabatic approximation suggested in Ref. [8] (see also Ref. [36]) and analogous to that used in the theory of nonstationary cavities and smoothly inhomogeneous waveguides (e.g., see Ref. [37]). According to this approach the electromagnetic fields are written in the quasimonochromatic form

$$\begin{Bmatrix} \boldsymbol{\mathcal{E}} \\ \boldsymbol{\mathcal{B}} \end{Bmatrix} = \frac{1}{2} \times \begin{Bmatrix} \boldsymbol{E}(y,t) \\ \boldsymbol{B}(y,t) \end{Bmatrix} \times \exp[i\varphi(t) - ih_0 x] + \text{c.c.} \tag{4}$$

with slow time-varying frequency $\omega(t) = d\varphi/dt$ and complex amplitudes $\boldsymbol{E},\boldsymbol{B}$ (c.c. means complex conjugate). The longitudinal wave number h_0 is fixed because of spatial homogeneity of the created plasma structure in the x direction. Indeed, averaged over the wave period $2\pi/\omega$ plasma density $N = \langle n \rangle$ depends, evidently, only on time t and transverse coordinate y, i.e., $N = N(y,t)$. The wave mode (4) develops in time continuously ''adiabatically'' following the time variation of the plasma density N. The amplitude and frequency may change significantly after a considerable length of time (much longer than $2\pi/\omega$). For every moment of time the frequency and transverse structure of the wave mode is found from the solution of an eigenvalue problem on the basis of the stationary wave equation for the magnetic field B_z and relations for the electric field components:

$$\varepsilon \frac{\partial}{\partial y}\left(\frac{1}{\varepsilon} \frac{\partial B_z}{\partial y} \right) + \left(\frac{\omega^2}{c^2}\varepsilon - h_0^2 \right) B_z = 0, \tag{5a}$$

$$\omega \varepsilon E_y = c h_0 B_z, \tag{5b}$$

$$\omega \varepsilon E_x = -ic \frac{\partial B_z}{\partial y}, \tag{5c}$$

where $\varepsilon = 1 - \omega_p^2/\omega(\omega - i\nu)$ is the complex permittivity of the plasma, ω_p and ν are the plasma frequency and electron collision frequency, respectively. The slowly varying in time transverse spatial distribution of the plasma frequency $\omega_p(y,t) = [4\pi e^2 N(y,t)/m]^{1/2}$ is determined by evolution of the average plasma density $N(y,t)$. Equation for $N(y,t)$ follows from Eq. (2) and has the form

$$\frac{\partial N}{\partial t} = 4\Omega (N_g - N) \sqrt{\frac{3}{\pi} \frac{\mathcal{E}_a}{|\boldsymbol{E}|}} \exp\left(-\frac{2}{3} \frac{\mathcal{E}_a}{|\boldsymbol{E}|} \right). \tag{6}$$

With the fixed dependence along the x direction, the solution of Eq. (5a) determines, at each instant of time t [for given distribution of the plasma density $N(y,t)$], the eigenfunction $B_z(y,t)$ and eigenvalue $\omega(t)$. Undefined time-dependent normalization factor in the eigenfunction $B_z(y,t)$ is found from the equation for evolution of wave energy

$$\frac{\partial}{\partial t} \int_0^L \frac{|\boldsymbol{E}|^2}{8\pi} dy = -\int_0^L Q\, dy, \tag{7a}$$

$$Q = \left\langle \frac{e^2}{2m\omega^4} \frac{\partial n}{\partial t} \left(\frac{\partial \mathcal{E}}{\partial t} \right)^2 \right\rangle + \frac{|\boldsymbol{E}|^2}{8\pi} \frac{\nu \omega_p^2}{\omega^2}. \tag{7b}$$

This equation, first derived in Ref. [8], is a generalization of the intensity transport equation [18,19] for the case when transverse spatial modulations of the electromagnetic field and the plasma density occur. The equation (7) follows from the Poynting's theorem for the quasimonochromatic field and the equation for the electron current $\boldsymbol{j}$ in a time-varying plasma [19]:

$$\frac{\partial \boldsymbol{j}}{\partial t} + \nu \boldsymbol{j} = \frac{e^2 n}{m} \boldsymbol{\mathcal{E}}. \tag{8}$$

Physically, the left-hand side of Eq. (7a) represents a time derivative of wave energy per unit area of x,z plane; the right-hand side of this equation determines the energy losses and it is a sum of two terms [see Eq. (7b)]. The first term describes energy transfer to newly born electrons and should be calculated with taking into account fast variations of electron density. The second term describes electron collision losses. Direct energy losses by the electrons detachment are negligibly small in the validity region of the static formula (2) ($I \ll W_q$). The small collision term was kept in Eqs. (5) to avoid divergence of electric field in the point of plasma resonance and describe correctly dispersive electrodynamical properties of the system (the latter will be clarified in Sec. IV).

Set of Eqs. (5)–(7) is complete and will be used in the next section for numerical simulation of the evolving discharge.

III. RESULTS OF NUMERICAL SIMULATION

The concrete character of nonlinear temporal evolution of the gas discharge, described by the set of Eqs. (5)–(7), depends significantly on the values of dimensionless parameters $\mathcal{E}_0/\mathcal{E}_a$, $\mu = m\omega_0^2/(4\pi e^2 N_g)$, and $ch_0/\omega_0 = \cos\theta_0$. Parameter μ defines the ratio between the critical plasma density for initial electromagnetic field $N_{c0} = m\omega_0^2/4\pi e^2$ and possible maximum value of the created plasma density N_g.

Before presenting results of the simulation let us discuss the general pattern of the discharge evolution. In the beginning of the process, ionization occurs practically only in thin (compared to the wavelength) layers at the interference maximums of the standing wave's electric field (i.e., near $y = 0, \pm 2L, \ldots$). With growth of the plasma density in the layer due to ionization, the component of electric field E_y increases inside the layer as well according to the quasistatic relation $E_y = \text{const}/\varepsilon$. The last relation follows from the continuity of the electric displacement vector across the thin layer (which is similar to the case of a plane capacitor). Note that the wave frequency changes insignificantly at the initial stage of the process because properties of the medium vary little and only in the small part of the periodicity interval 0

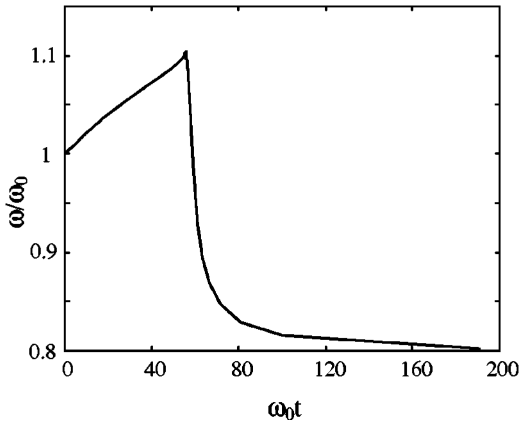

FIG. 2. Temporal evolution of the electromagnetic field frequency. Sharp decrease of frequency occurs at $\omega_0 t \approx 56$.

$<y<L$. So the variation of ε in the plasma layer is caused mainly by the increase of ω_p and, therefore, ε decreases. The increase of E_y in the layer results in a growth of the ionization rate $\langle w \rangle$ and a further growth of plasma density. This leads to the growth of the E_y inside the layer. Thus, the effect of spatial sharpenning of the electric field occurs. The rate of ionization increases when the real part of ε becomes closer to zero and it reaches peak value when $\mathrm{Re}\varepsilon = 0$, i.e., plasma density goes through the critical value and plasma resonance occurs. After that the growth of plasma density in the slab slows down and the process of slab widening begins to dominate. The described process has common features with the process of the plasma-resonance ionization instability considered in Ref. [8].

The main condition of the applicability of the adiabatic approximation used here is the slowness of temporal variations of the field and plasma density. By choice of the parameters $\mathcal{E}_0/\mathcal{E}_a$, μ, and ch_0/ω_0 it is possible to satisfy this condition and, at the same time, to provide significant sharpenning of the initial electric field distribution in the course of the discharge evolution. This case is realized, for example, for the following set of parameters, used in the numerical simulations: $ch_0/\omega_0 = 0.95$ ($\theta_0 \approx 18.2°$), $\mathcal{E}_0/\mathcal{E}_a = 0.025$, and

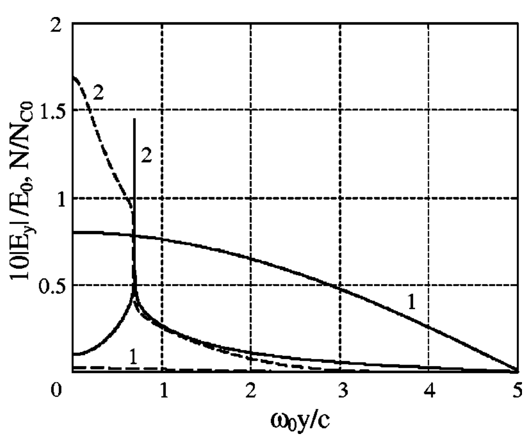

FIG. 3. Transverse structure of the electric field amplitude $|E_y|$ (solid line) and plasma density N (dashed line) for two instants of time: $\omega_0 t = 2$ (1); 190.72 (2).

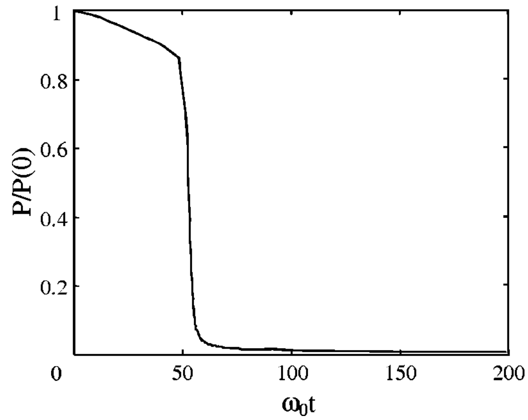

FIG. 4. Temporal evolution of wave's power flow P across the plane y,z (per unit interval of z and $0<y<L$) normalized to the initial value $P(0) = (c/4\pi)B_0^2 L \cos\theta_0$. Sharp decrease of $P/P(0)$ at $\omega_0 t \approx 56$ corresponds to the frequency drop in Fig. 2.

$\mu = 0.4$. Relative collission frequency ν/ω_0 was taken equal to 0.3 and $\Omega/\omega_0 = 50$. The assumed value of ν/ω_0 is in a good agreement with estimations for real experimental situations. It may correspond, for example, to the ionization of H_2 gas by the laser pulses of wavelength band $1-10$ μm; corresponding gas pressure is 60-0.6 atm, respectively.

Results of the computer simulation are presented in Figs. 2–5.

Temporal evolution of the electromagnetic field frequency, which occurs simultaneously with the discharge evolution, is shown in Fig. 2. A distinguishing feature of the graph is a sharp decrease of the frequency below the initial value ω_0 at the moment when $\omega_0 t \approx 56$. This sharp frequency decrease comes immediately after its slow growth that initially takes place. For some values of the parameters, calculations give frequency drop up to $\omega_{min}/\omega_0 \approx 0.7$. Such behavior of the electromagnetic field frequency is extraordinary for the systems with time-growing plasma density, in which frequency upshifting of radiation is usually discussed [11]. The difference between the results presented here and in pre-

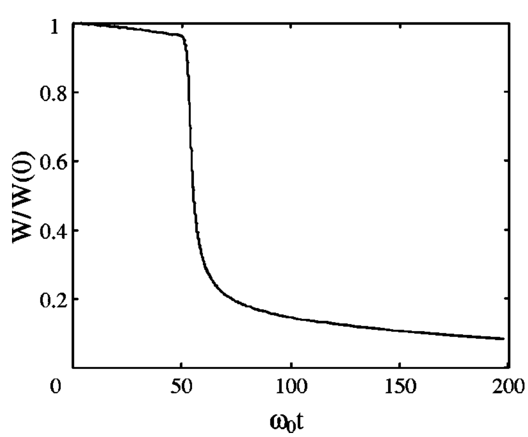

FIG. 5. Temporal evolution of the energy of the wave W (per unit area of x, z plane and $0<y<L$) normalized to the initial value $W(0) = B_0^2 L/4\pi$. Sharp decrease of $W/W(0)$ at $\omega_0 t \approx 56$ corresponds to the frequency drop in Fig. 2.

vious works is explained by the fact that earlier only plane waves or wide beams in a homogeneous and quasihomogeneous time-varying plasma were considered, whereas in our case the dramatic changes in the transverse structure of the wave and the medium occur. These changes are demonstrated in Fig. 3 where transverse distributions of the electric field $|E_y|$ and plasma density N for two instants of time (before and after $\omega_0 t \approx 56$) are presented. It follows from Fig. 3 that at the definite stage of the discharge evolution ($\omega_0 t \approx 56$) the initial bulk (standing) wave converts into surface (slow) wave localized at the interface between the dense and rarefied plasmas. Saying successively, within every spatial interval of the field periodicity the bulk wave creates plasma waveguide—the layer of an overdense plasma—that in its turn traps the wave via its conversion into localized mode of the waveguide. In other words, the self-sustained plasma waveguide arises and, thus, the self-channeling of the electromagnetic radiation is realized.

Figure 4 shows the time dependence of the wave's power flow $P(t)$. The drop of $P(t)$ at $\omega_0 t \approx 56$ is provided by both the slowing down of the wave and sharp decrease of the wave energy $W(t)$ (see Fig. 5). $W(t)$ decreases because ionization rate reaches peak when density of the time-growing plasma goes through the critical value and plasma resonance occurs. Peaking of the ionization rate gives peaking of the energy transfer from the wave to newly born electrons.

It also follows from Fig. 2 that after the sharp drop the frequency continues to decrease but slightly and slowly. The slow decrease of the frequency is accompanied by the decrease of the phase velocity, power flow, and energy of the guided wave (see Figs. 4 and 5).

IV. THEORETICAL MODEL

To reveal the underlying physics of the self-downshifting and bulk-to-surface-wave conversion phenomena we attempt here to construct a simple model of the field-plasma interaction. First of all, let us note that the development of the discharge leads to formation of the spatial distribution of the plasma density with a sharp step between two practically homogeneous regions of dense and rarefied plasma (see Fig. 3). This fact allows us to model this nonstationary distribution by a double-layer (on every spatial period L) structure consisting of a homogeneous plasma slab of density $N(t)$ and thickness $d(t)$ and a vacuum layer of thickness $L-d$. The parameters of the plasma slab $N(t)$ and $d(t)$ are slow (compared to the wave period $2\pi/\omega$) functions of time those evolve according to described in Sec. III general pattern of the discharge evolution. Namely, we assume that, at first, the plasma density in the layer $N(t)$ grows in time at fixed value of d. Then, when the plasma density reaches a value that is slightly greater than the critical plasma density, the growth stops and the thickness of the layer $d(t)$ begins to increase. The slowness of the parameters variations allows us to consider the evolution of a wave, which propagates along the spatially periodic structure, adiabatically, i.e., to write the wave in the form as given by Eq. (4). Within the adiabatic approximation, the time dependence of the wave frequency $\omega(t)$ may be found from the dispersion equation $D(\omega,h)=0$, which corresponds to the instant state of the waveguiding structure. The longitudinal wave number in the dispersion equation must be fixed equal to h_0: $h=h_0$. For given N and d the dispersion equation is easily derived from Eqs. (5) and usual boundary conditions of continuity of E_x and B_z at the plasma-vacuum boundary. Finally, we arrive at the equation of the form

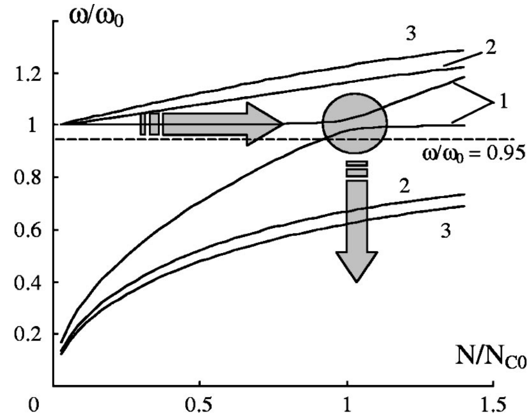

FIG. 6. Frequencies of the two lowest modes of the model structure versus the density of the plasma slab N normalized to the critical (for the wave of frequency ω_0) value N_{c0} for different values of the slab thickness: $d/L=0.001$ (1); 0.141 (2); 0.281 (3). Segments of the dispersion curves lying below the level $\omega/\omega_0=0.95$ (marked by the dashed line) correspond to slow surface wave.

$$\frac{\kappa_p}{\kappa_v}\frac{\exp(\kappa_p d/2)-\exp(-\kappa_p d/2)}{\exp(\kappa_p d/2)+\exp(-\kappa_p d/2)} - \varepsilon\frac{\exp[\kappa_v(L-d/2)]+\exp[-\kappa_v(L-d/2)]}{\exp[\kappa_v(L-d/2)]-\exp[-\kappa_v(L-d/2)]}=0, \tag{9}$$

where ε is the dielectric function of the plasma slab and $\kappa_v=\sqrt{h_0^2-\omega^2/c^2}$, $\kappa_p=\sqrt{h_0^2-\varepsilon\omega^2/c^2}$ are the transverse wave numbers in vacuum and in the plasma, respectively. Equation (9) is transcendental and, therefore, the dependence $\omega(t)$, initiated by variations of $N(t)$ and $d(t)$, cannot be expressed in explicit form. Moreover, Eq. (9) does not have a unique solution. There is a set of branches corresponding to various modes of the periodic plasma-vacuum structure. We investigated Eq. (9) numerically for the same value of parameter ch_0/ω_0 as in Sec. III, i.e., for $ch_0/\omega_0=0.95$. Note that initial frequency ω_0 appears in Eq. (9) through the parameter $L=\pi/2g_0=\pi c/(2\omega_0 \sin\theta_0)$ that corresponds to the original nonlinear problem formulated in Sec. II. At first, we neglect collissions. Figure 6 shows the frequencies of the two lowest modes of the structure in dependence on the plasma density N for different values of the plasma slab thickness d. The peculiarity of the graph is that for small thickness of the slab ($d/L=0.001$) the narrow interval of plasma density near $N/N_{c0}=1$ exists where the upper and lower branches of the curve $\omega(N)$ (curve 1 in Fig. 6) are very close to each other (this region in Fig. 6 is marked by

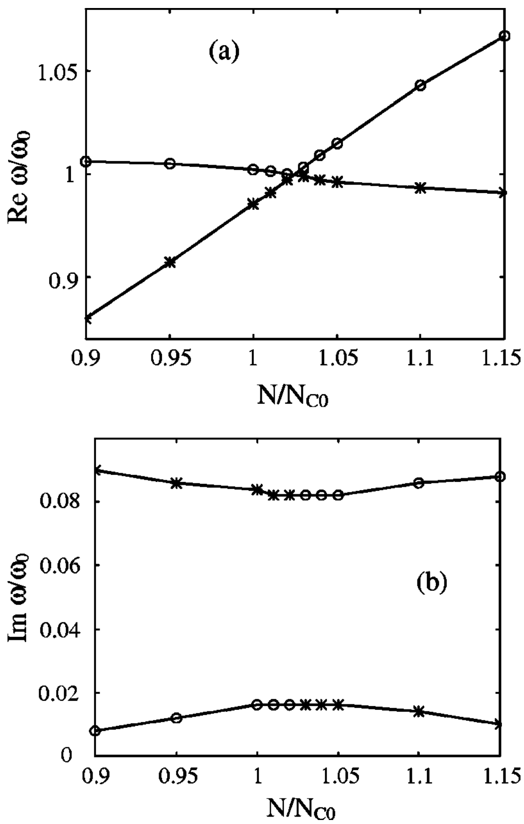

FIG. 7. Real (a) and imaginary (b) parts of frequencies of the model structure modes versus the plasma density N in the case when collisions are included. The plasma slab thickness is $d/L = 0.003$, collision frequency is $\nu/\omega_0 = 0.2$. Circles and stars mark the segments attributed to different modes.

the shaded circle). The calculations show that the branches touch each other at the point $N/N_{c0} = 1$ in the limit $d \to 0$. When parameter d/L increases, the upper and lower branches diverge, i.e., they move in the graph into opposite directions (curves 2 and 3 in Fig. 6). Thus, when collissions are neglected the modes of the structure cannot converge one into another adiabatically for any finite thickness d. The principal point of the consideration is in the fact that for small (but not vanishing) values of d/L the upper and lower branches merge together when collisions are taken into account. It is demonstrated in Fig. 7(a) that shows Reω as a function of plasma density N. Another important point is that adiabatic evolution of the mode through the point where the branches intercross should be accompanied by the transition from one branch to another. It is dictated by the physically natural conditions that imaginary part of frequency Imω should change continuously with N and the real part of frequency Reω should change smoothly (without the jump of derivative). It follows clearly from Fig. 7(b) that continuity of Imω leads to transition from the upper branch to the lower one and vice versa [see also Fig. 7(a)]. By using Figs. 6 and 7, we can trace the initial wave evolution with the variations of parameters $N(t)$ and $d(t)$. First of all, we note that when $N = 0$ the lower branch in Fig. 6 equals zero ($\omega = 0$) whereas the upper one equals ω_0. Evidently, it is the upper branch that corresponds to the initial wave at $N = 0$. According to the above described scenario of the discharge evolution, the plasma density grows initially in a thin layer near the maximum of the electric field amplitude. Correspondingly, the representing point in Fig. 6 moves along the upper branch of curve 1 to the right (it is marked by arrow in Fig. 6). The wave frequency practically does not change on this stage of evolution. Further, when the representing point reaches the region of the modes interaction (the shaded circle in Fig. 6), i.e., the plasma density reaches the critical value, a transition from the upper branch to the lower one takes place. Simultaneously, the plasma density stops to grow and the plasma slab thickness starts to increase. The representing point in Fig. 6 moves down from curve 1 to the lower branches of curves 2, 3, and so on (this process is marked by the vertical arrow). On this stage of the discharge evolution significant frequency downshifting occurs. It is worth noting that when the representing point falls below the level $\omega/\omega_0 = 0.95$ in Fig. 6 it means qualitative change of the wave type: bulk wave converges into surface one.

V. CONCLUSIONS

To conclude, we have presented a nonlinear phenomenon of bulk-to-surface-wave self-conversion for intense laser radiation with inhomogeneous transverse structure in an ionized gas. This phenomenon consists in the creation of a layered plasma structure by the electromagnetic field and the conversion of the radiation into surface waves that are guided by the structure and that support in their turn ionization in the layers. The dynamic process of the formation of such a self-sustained plasma waveguiding structure is traced numerically for two intersecting plane waves in a gas. Tunnel mechanism of ionization is considered. Parameters for numerical simulation were taken to provide adiabatic character of the laser-produced plasma development and bulk-to-surface-wave conversion process. The process is shown to be accompanied by significant frequency downshifting of the electromagnetic radiation. The underlying physical mechanism has also been explained on the basis of a simple analytical model of periodic structure with homogeneous plasma layers.

To focus on the principal features of the phenomenon, we considered the initial value problem for two intersecting waves of infinite transverse size (plane waves). Certainly, for a description of the real experimental situation with two crossed laser beams (see, e.g., Ref. [38]) or a single beam with a stray quasiperiodic transverse modulation a more complicated initial-boundary problem should be addressed. This problem should include the limited transverse size of the laser field distribution, longitudinal inhomogeneity of laser pulse envelope, process of entering the radiation into gas, and processes behind the intersection region of two beams or focal region for a one beam. Our consideration was aimed to demonstrate the principal possibility of the bulk-to-surface-wave self-conversion phenomenon on a simplest model.

Although the results obtained concern directly the realization of self-channeling in the intersection region, the presented phenomenon of bulk-to-surface-wave self-conversion may be of interest as a possible mechanism of self-channeling of intense laser pulses beyond the intersection region. Indeed, we can await that the self-sustained plasma waveguiding structure shown to arise in the intersection region can prolong itself beyond this region due to ionization of the gas by the fields of surface waves propagating in forward direction along the plasma layers of the structure. Evidently, beyond the intersection region the self-sustained waveguiding structure will fade with x due to depletion of the surface waves energy.

One can conclude from our analysis that possibility of the bulk-to-surface-wave self-conversion phenomenon is relied upon the existence of the dispersion curve of the created plasma structure whose different segments correspond to waves of different (bulk and surface) types (see lower curve 1 in Fig. 6). However, it is not peculiarity of strictly periodic structures only. It is easy to show, for example, that there is a similar dispersion curve in a three-layer structure ''plasma-vacuum-plasma.'' This fact allows us to await that the bulk-to-surface-wave self-conversion phenomenon can be realized even for intersection of laser beams under small angles when only a few spatial oscillations of electric field occur across the intersection region.

In the above considered adiabatic conversion process electron collisions play significant role: the relative collision frequency should not be too small [see Figs. 6 and 7(a,b)]. It means that the plasma density and, therefore, gas pressure should be high enough. Thus, one can await that the phenomenon of bulk-to-surface-wave self-conversion may be even more significant for dynamics of laser pulses interaction with solid and liquid media. Also, it would be interesting to consider similar effect for faster ionization processes where adiabatic approximation breaks down.

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