Electrodynamic mechanisms that limit the electron concentration in a laser spark

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The joint evolution of the electromagnetic field and of the plasma in a nonequilibrium discharge produced in a gas by a focused laser beam is investigated. The maximum electron density in the breakdown wave $N_{e\max}$, which is determined by the electrodynamic mechanisms that limit the electron avalanche (by refraction and absorption) is obtained. The value of $N_{e\max}$ is proportional to the square of the convergence angle of the beam and increases (logarithmically) with decrease in the initial (pre-breakdown) density.

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An important problem in the theory of a laser spark, just as of any other high-frequency discharge, is the determination of the electron-density level that limits the avalanche breakdown process. In the earlier experimental and theoretical investigations (see the monograph[1] and the literature cited therein) there was realized and analyzed the situation that is the simplest from the point of view of the indicated problem, wherein the growth of the electron density continues until the gas is completely ionized. At the present time, however, (principally in connection with the planned expansion of the research on the laser-spark to include longer wavelengths,[2,3] no less interest attaches to the study of the dynamics of the optical (or quasioptical) discharge under conditions when the "electrodynamic" mechanisms that stop the avalanche become substantial; these mechanisms are connected with the weakening of the field of the wave in the plasma by absorption or refraction. In the present paper we consider a discharge in the field of a single-mode long-focus beam, in which the indicated mechanisms limit the electron density to a rather low level, and by the same token prevent effective heating of the gas and stop the development of the discharge during the initial nonisothermal stage. The problems of the kinetics of the elementary process in such a discharge are in general the same as for the start of the avalanche process, and can be regarded in the main as solved; the principal focus of attention in the problem is now on the electrodynamics, i.e., on the investigation of the joint evolution of the wave field and of the plasma at a certain albeit quite restricted type of nonlinearity.

1. In the electron balance equation

$$\frac{\partial N_e}{\partial t}=\gamma N_e,\qquad N_e=N_{e0}\exp\int_0^t\gamma\,dt \tag{1}$$

we regard the rate γ of the avalanche (the difference between the frequencies of the ionization and of the electron loss) as a known (increasing) function of the amplitude of the electric field, $\gamma=\gamma(E)$, and assume that the unperturbed beam is stationary (it is turned on instantaneously at the instant $t=0$) and Gaussian:

$$E_0(r,z)=E_F\left(1+\frac{z^2}{l_F^2}\right)^{-1/2}\exp\left(-\frac{r^2}{2a^2}\right). \tag{2}$$

Here z is the axial coordinate measured upstream from the focal plane along the flux of the incident radiation, r is the distance from the beam axis, $a^2=a_F^2+\theta^2z^2$; and $l_F=ka_F^2$ are the transverse and longitudinal dimensions of the focal region, $\theta=(ka_F)^{-1}$ is the beam focusing angle (it is assumed that $\theta\ll1$), $k=2\pi/\lambda=\omega/c$, ω is the field frequency, and E_F is the amplitude of the field at the center of the focal spot.

At any section of the beam where z is constant (where $\gamma>0$) the growth of the electron density N_e continues until refraction and absorption of the wave weaken substantially the field in the plasma produced ahead of this section (in the region of larger z). As a result, on the leading front of the breakdown wave,[1)] which propagates in the region $z<0$ in the direction of decreasing unperturbed avalanche velocity $\gamma_0=\gamma(E_0)$ (i.e., in a direction opposite to the incident radiation), mutually compatible drops take place in the field amplitude E and in the electron density N_e, whereas in the region $z<0$, as everywhere behind the front of the breakdown wave, the ionization stops rapidly. To estimate the maximum density $N_{e\max}(z)$ reached at the instant t_m of the strong slowing down of the avalanche, it is necessary to determine, at $t=t_m$, the characteristic structural parameters of the screening region adjacent to the given cross section z. These parameters are the distance Δz along the axis over which the argument of the exponential in (1) decreases by unity, and the width b of the radial profile (which is approximated in the paraxial region by the parabola $N_e(r)\sim1-r^2/b^2$). At sufficiently small values of the initial ("bare") density $N_{e0}(\ln(N_{e\max}/N_{e0})\gg1)$, when the principal part of the growth of N_e takes place in the unperturbed field E_0 at a constant rate $\gamma_0=\gamma(E_0)$, the argument of the exponential in (1) for $T=t_m$ can be approximately set equal to $\gamma_0(z)t_m$, so that we can obtain for Δz and b the following expressions:

$$\Delta z\approx\begin{cases}l_F(2/\beta Q)^{1/2}, & z<l_F(2\beta Q)^{-1/2}\\ a^2/\beta Q\theta^2z, & z>l_F(2\beta Q)^{-1/2}\end{cases}, \tag{3}$$

$$b\approx a(2/\beta Q)^{1/2}, \tag{4}$$

where $Q=\ln(N_{e\max}/N_{e0})\gg1$ (as a rule, $Q\approx30-40$); $\beta=(E/(E/\gamma)d\gamma/dE$. In the case of rapid ionization of the excited atoms by the incident radiation we have $\beta=2(\gamma\propto E^2)$. If no such ionization takes place and the loss of electron energy to excitation is large, the value of β is 2–3 times larger.[1,6] Under the condition

$$ka_F\gg(\beta Q/2)^{1/2}\quad((\beta Q/2)^{1/2}\approx5-7), \tag{5}$$

which is more stringent than the simple condition

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$ka_F = \theta^{-1} \gg 1$ that the beam be paraxial, the size of the discharge is large compared with the wavelength and its action in the field can be described within the framework of quasi-optical concepts.

2. We estimate now the maximum electron density on the beam axis, taking into account separately the action of the two aforementioned factors of quasioptical weakening of the field-absorption and refraction—and regarding the plasma of the discharge as a dielectric with a complex dielectric constant

$$\varepsilon = 1 - n - i\frac{\nu}{\omega}n, \tag{6}$$

where ν is the effective electron collision frequency, $n = N_e/[N_c(1+\nu^2/\omega^2)]$ is the dimensionless density, and $N_c = m\omega^2/4\pi e^2$ is the critical density in the absence of collisions. As we shall see, when the condition (5) is satisfied the value of $|\varepsilon|$ is close to unity ($n(1+\nu/\omega \ll 1$), so that we can neglect the direct reflection of the wave from the discharge and the large-angle scattering.

When account is taken of absorption (without refraction), the field on the beam axis takes the form

$$E(0,z) = E_0(0,z)\exp\left(-\frac{\nu}{2c}\int_{-\infty} n(z)\,dz\right). \tag{7}$$

We estimate the maximum concentration as that value of n at which the avalanche velocity is strongly decreased $\delta\gamma/\gamma \approx (\delta E/E)\beta \approx -1$, i.e., the argument of the exponential in (7) reaches a value of the order of $1/\beta$:

$$\frac{\nu}{2c}\int_{-\infty}^{z} n\,dz \approx \frac{\nu}{2c} n_{\max}\,\Delta z \approx 1/\beta \tag{8}$$

whence, with allowance for (3), we get

$$n_{max}(z) = \begin{cases} (2Q/\beta)^{1/2}\dfrac{\omega}{\nu}\theta^2, & z < \dfrac{l_F}{(2\beta Q)^{1/2}} \\ 2Q\dfrac{\omega}{\nu}\theta^2\dfrac{z/l_F}{1+(z/l_F)^2}, & z > \dfrac{l_F}{(2\beta Q)^{1/2}} \end{cases}. \tag{9}$$

The largest value of $n_{\max}(z)$ is reached at $z = l_F$:

$$n_{\max}(l_F) \approx \frac{\omega}{\nu}\theta^2 Q. \tag{10}$$

The refraction in the paraxial region of the beam in the absence of absorption (as $\nu/\omega \to 0$) can be taken into account with the aid of the equation for the radius of the ray tube $r(z)$:

$$\frac{d^2r}{dz^2} = \frac{nr}{b^2} + \frac{1}{2k^2}\frac{\partial}{\partial r}\left(\frac{\Delta_\perp E}{E}\right). \tag{11}$$

The amplitude $E(z)$ of the field on the axis is connected with the radius $r(z)$ by the relation $Er = \text{const}$. At small n, the relative broadening of the ray tube $\delta r/r$ due to the refraction and the decrease of the amplitude $\delta E/E = -\delta r/r$ can be estimated on the basis of (11) by a perturbation method, using the expressions (3) and (4) given above for the characteristic scales Δz and b of the density variation.

Since the width of the region occupied by the plasma is small compared with the width of the beam ($b \ll a$), in the estimate of $\delta r/r$ in (11) it is necessary, generally speaking, to take into account also the correction to the diffraction term (the second term in the right-hand side). In the focal region of the beam ($z \gg l_F$), where the wave parameter $\Delta z/kb^2 \gg 1$, the diffraction prevents the appearance of narrow and deep field dips in the paraxial region, which might occur when account is taken of only the refraction correction. In the region $z \gg l_F$ the parameter $\Delta z/kb^2 \ll 1$ and the diffraction are inessential. In the general case, representing the field in the form $E(r) = E_0(r) + \delta E(r)$ and recognizing that at small n the correction $\delta E(r)$ is localized in a region of width $b \ll a$, we obtain from (11) the following estimate for the relative broadening of a ray tube of radius $r \lesssim b$

$$\frac{\delta r}{r} \approx n\left(\frac{\Delta z}{b}\right)^2\left[1+\left(\frac{\Delta z}{kb^2}\right)^2\right]^{-1}. \tag{12}$$

The maximum density $n_{\max}$, just as in the preceding case, is obtained from the condition that the avalanche velocity decrease substantially $\delta\gamma/\gamma \approx \beta\delta E/E \approx -1$. Putting $\delta r/r \approx -\delta E/E \approx 1/\beta$ in (12) and taking (3) into account, we obtain for both values of z indicated in (3)

$$n_{\max} \approx \theta^2 Q, \quad N_{e\,\max} \approx N_c\theta^2 Q. \tag{13}$$

It is obvious that the major of the two considered field-weakening mechanism is the one that stops the avalanche at a lower electron density, i.e., that yields lower values of $n_{\max}$. Comparing the obtained expressions for $n_{\max}(z)$ we arrive at the following results. Under the condition $\nu/\omega \ll 1$, which is satisfied for most optical-breakdown experiments, the principal role in the basic breakdown zone is played by refraction, i.e., the maximum density should be estimated from Eq. (13). Only at a sufficiently large distance from the focus $z < l_F\omega/\nu$ (if the breakdown wave reaches this region), and under the condition $(2/\beta Q)^{1/2} < \nu/\omega \ll 1$ the absorption predominates also in a small vicinity of the focus $z < l_F(2\beta Q)^{-1/2}$. At $\nu/\omega \gg 1$ (in the high-pressure region) the absorption predominates everywhere and the density is determined by (9); at the point $z = l_F$ we have

$$N_{e\,\max}(l_F) \approx N_c\frac{\nu}{\omega}\theta^2 Q. \tag{14}$$

Thus, the electron density in the breakdown wave is proportional to the square of the focusing angle θ (i.e., it is inversely proportional to the area of the focal spot)[2)] and increases (logarithmically) with increasing initial density N_{e0}. The last circumstance, which is somewhat paradoxical at first glance, can be explained in the following manner: the smaller N_{e0}, the longer the time $t = \gamma^{-1}\ln(N_e/N_{e0})$ needed to increase N_e to any specified level, the larger the density differentials produced during that time in any specified segment of the z axis, and consequently the smaller the characteristic scale Δz of the growth of N_e in the breakdown wave and the larger should be $N_{e\,\max}$ in order for the weakening of the field over this scale to become noticeable and to be able to stop the avalanche.

3. The breakdown-wave front, whose duration is $\tau \approx \gamma^{-1}$ and which moves with velocity $v(z) \approx \gamma\Delta z$, is obviously the zone of the more active energy exchange between the radiation and the discharge—the density here is of the order of the maximal one, and the field has not yet been strongly enough decreased compared with the unperturbed one. The fore-going analysis is valid under the condition that no considerable heating of the gas takes place in this zone (the temperature increment $\Delta T_m < 10^3$K) and no high degree of ionization is reached ($N_{e\,\max} < N_m$, N_m is the density of the neutral particles).

The increase ΔT_m of the gas temperature depends on the degree of ionization $N_{e\,max}/N_m$ of the average energy of the electrons T_e and on the time γ^{-} of the passage of the active zone through the given point. The quantities T_e and γ, in turn, are determined by the ratio of the energy w of the oscillatory motion of the electrons in the wave field to the potential I of the ionization of the atoms. At the usual radiation intensities for pulsed laser breakdown, w lies in the range $\delta I < w < I$, where δ is the fraction of the energy lost by the electrons in the collisions. Under these conditions, at incomplete ionization of the gas $T_e \approx I \approx 10$ eV, $\tau = \gamma^{-1} \approx \nu^{-1} I/w$ and ΔT_m is quite small even if one includes in it the entire energy given up to the electrons in the elastic and inelastic collisions:

$$\Delta T_M \approx T_e \delta \nu \tau \frac{N_{e\,max}}{N_m} \approx I \frac{\delta I}{w} \frac{N_{e\,max}}{N_m}. \qquad (15)$$

Behind the breakdown-wave front the heating of the gas by the electrons continues, but it now occurs actually in the absence of a field and can be of importance only during the succeeding stages of the discharge.

We present in conclusion some quantitative estimates, assuming $\ln(N_{e\,max}/N_{e0}) = 30$, $I\delta/w = 10^{-1}$. For a neodymium laser ($N_c \approx 10^{21}$ cm^{-3}), at pressures $p \approx 1$ atm, even in the case of very small angles $\theta \approx 3 \times 10^{-2}$, the considered breakdown-wave regime does not occur in a cold gas ($N_{e\,max} \approx N_m$; T_e, $T_m > I$). For a CO_2 laser ($N_c \approx 10^{19}$ cm^{-3}) at $p = 1$ atm and $\theta = 3 \times 10^{-2}$ we have $N_{e\,max} \approx 3 \cdot 10^{17}$ cm^{-3} and $\Delta T_m \approx 100$ K. In the region of longer wavelengths and at not too low pressures, there is practically no heating of the gas: at $\lambda = 3 \cdot 10^{-2}$ cm ($N_e \approx 10^{16}$ cm^{-3}), $p = 0.1$ atm, and $\theta = 10^{-1}$ we have $N_{e\,max} \approx 3 \cdot 10^{15}$ cm^{-3} and $\Delta T_m \approx 10$ K.

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[1)]The kinematics of the breakdown wave in a given field (without allowance for the screening action of the produced plasma) was investigated in Refs. 4 and 5.

[2)]A similar result was obtained in an investigation of the stationary regime of a nonequilibrium discharge in the field of two waves that converge at a small angle,[7] as well as in a numerical simulation of the dynamics of a discharge in a wave beam.[8]

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Probe investigations of electric fields produced in air near a laser spark

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We registered two components of the electric field produced near the plasma due to breakdown of air by CO_2 laser radiation. The appearance of a rapidly alternating component with duration of the order of the duration of the laser pulse is connected with the separation of charges on the plasma front that propagates in a direction opposite to that of the laser beam. It is shown that the cause of the slowly varying component of the field is the photoeffect, induced by the plasma radiation, in the gas on the surfaces of the bodies surrounding the probe.

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In our preceding studies[1,2] we investigated signals from a differentiating electric probe located near air-breakdown plasma initiated on the surface of a conducting target. It was shown that registered signals were due not to radiation from the plasma, but to separation of the charges in the moving plasma front. The duration of these signals did not exceed the duration of the laser pulse. In the present paper we report probe measurements of the potentials φ of the electric fields near the plasma of air breakdown by radiation of a pulsed CO_2 laser on a dielectric target and in the absence of a target.

It is shown that if a "transmitting" rather than a differentiating electric probe is used (producing a signal $\propto\varphi$) there can be registered near the plasma potentials due not only to the proper electric field of the plasma produced by the optical breakdown of the air, but also

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