Fast gas-ionization wave in a laser beam

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The threshold intensity of laser radiation corresponding to the appearance of a fast ionization wave is determined. The velocity of the nonstationary motion of the plasma front is calculated. The nonstationary motion is characterized by an increase of the front velocity with increase of the gas pressure and this distinguishes it from stationary regimes. The dependences of the front velocity and of plasma temperature behind the front on the laser beam radius are found for the stationary motion of the fast ionization wave. The results of the calculations are in good agreement with the experimental data.

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INTRODUCTION

Rapid superdetonation motion of an ionization front in a gas opposite to laser radiation, was investigated in Refs. 1–10. The intensity of the laser radiation in the experiment was much lower than the gas breakdown threshold q^*, so that the "priming" plasma was produced by a pulse of an auxiliary laser at the intersection of the beams,[2] or else the laser beam was aimed on a graphite surface and the motion of the fast ionization wave started from a spark on the surface.[3–5]

We consider the mechanism whereby the ionization front is transported in this regime.[6,7] The radiation of the spark produces a certain distribution of free electrons in the gas.[8] In a laser beam, these electrons initiate an avalanche: the number of free electrons per unit volume increases exponentially:

$$n(x,\ t)\sim n_0(x)\,e^{t/\tau}.$$

Here τ is the time constant of the avalanche.[8] The coordinate x is reckoned from the breakdown point in a direction opposite to the laser radiation. The number $n_0(x)$ of free electrons produced in a unit volume by the ionizing radiation of the spark, decreases with increasing distance from the point of breakdown, leading to a "phase shift" in the development of avalanches at different sections of the beam. This simultaneous but phase-delayed development of the electron avalanches on a segment of the laser beam manifests itself as motion of an ionization front over this segment and constitutes a fast ionization wave. If the laser pulse duration τ_L is long enough, the plasma front can move a considerable distance along the laser beam from the point of the initial breakdown. The priming photoionization of the cold gas is due in this case mainly not to the initial spark but to the leading plasma radiation emerging from behind the wave front. Moving away from the spark, the fast ionization wave reaches a stationary regime of propagation. It is these steady-state waves which were considered in Refs. 6 and 7 to explain the experimental results.[3–5]

The present article is devoted to a determination of the threshold of development of fast ionization wave, to calculation of velocity of the unsteady motion, and to calculation of the dependence of the front velocity and of the plasma temperature on the laser-beam radius.

THRESHOLD FOR THE ONSET OF FAST IONIZATION WAVE

We calculate the minimum laser-radiation intensity at which a fast ionization wave can still be detached from the priming plasma. We isolate within the laser beam a thin layer of gas perpendicular to the beam axis, and calculate the function $n(t)$ as the wave front passes through this layer. The first electrons appear as the result of photoionization of the atoms of the plasma radiation. The rate of photoionization of the gas on the laser beam axis is[6]

$$\dot{n}_{\rm ph}(x)=2\pi\int\limits_0^{\operatorname{arctg}(R_1/x)}\sin\vartheta\,d\vartheta\int\limits_{I_i/h}^{\infty}\frac{\varkappa_\nu^{\rm ph\,0}\,J_\nu(x,\vartheta)}{h\nu}\,d\nu. \tag{1}$$

Here $\varkappa_\nu^{\rm ph,0}$ is the spectral coefficient of bound-free absorption in a cold gas ahead of the wave, J_ν is the plasma radiation intensity, ϑ is the angle between the laser-beam axis and the vector direction for the point x toward the radiating volume, and R_1 is the transverse dimension of the radiating volume (the laser-spark radius if the initial stage of the ionization wave motion is considered, or the radius R of the laser beam if the wave has reached the stationary propagation regime).

In the field of high-power laser radiation, the effective $\tilde{\nu}$ consists of the frequency ν_i of the direct ionization and the frequency $\nu_{\rm ex}$ of the level excitation, since the excited levels are ionized rapidly by radiation or by electron impact.[9,6]

Taking electron diffusion into account, we write down the ionization-kinetics equation in the form

$$\dot{n}=\dot{n}_{\rm ph}+\tilde{\nu}n+D_e\Delta n-\dot{n}_-. \tag{2}$$

Here D_e is the coefficient of the free diffusion of the electrons and $\dot{n}_-$ is the recombination rate. The electron recombination stops the exponential growth of the function $n(t)$ as the equilibrium value $n_{\rm eq}(T_e)$ is approached, but the most complicated and slowest for the avalanche development are the earlier stages of electron multiplication,[9,10] for it is they that determine the threshold intensity of the laser radiation; recombination can be neglected when the threshold is calculated. We neglect also the sticking of the electrons, since inert gases are usually used in the experiment. Introducing the diffusion length L, we can replace Δn by $-n/L^2$ and write down the solution of (2) in the form

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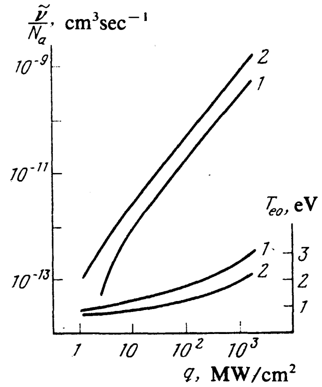

FIG. 1. Dependence of the electron temperature in an avalanche and of the frequency of the inelastic collisions of the electrons with the atoms (per atom) on the CO_2 laser radiation. 1) Argon; 2) xenon.

$$n(t)=\dot{n}_{ph}\tau(e^{t/\tau}-1), \tag{3}$$

where the avalanche time constant is

$$\tau=(\tilde{\nu}-D_eL^{-2})^{-1}. \tag{4}$$

The temperature of the electrons in the avalanche $T_{e0}(q)$ is determined by the equilibrium between the absorption of the laser radiation and the energy lost in the collisions:

$$\frac{4\pi e^2}{mc}\frac{n\nu_c}{\omega^2+\nu_c^2}q=\frac{2m}{M}n\nu_c\frac{3}{2}(T_{e0}-T)+n\tilde{\nu}\left(\bar{I}+\frac{3}{2}T_{e0}\right). \tag{5}$$

Here ω and q are the frequency and intensity of the laser radiation, $\nu_c(T_e)$ is the frequency of the elastic collisions of the electron, T is the temperature of the atoms and ions, $\bar{I}$ is the electron energy: $I_{ex}\leqslant\bar{I}\leqslant I_i$, and I_{ex} and I_i are the potentials of the first excited level and of the ionization of the atom. Substituting the explicit expressions for $\nu_c(T_e)$ and $\tilde{\nu}(T_e)$, we can determine the function $T_{e0}(q)$. By way of example, Fig. 1 shows results obtained for argon and xenon. The frequencies ν_{ex} and ν_i of the excitation and ionization of the atoms by electron impact were calculated in accordance with Refs. 10 and 11. The effect of frequency of the elastic collisions of the electron with the atoms was calculated by integrating the measured cross sections $\sigma_c(\varepsilon_e)$ (Ref. 12). The integration results are given in Ref. 7 for $1\leqslant T_e\leqslant 6$ eV. The effective frequency of the elastic collisions between the electron and the ions was calculated from known expressions.[13] The plots shown in Fig. 1 correspond to the experimental conditions of Refs. 3–5, namely CO_2 laser, $R\approx 1$ cm, gas pressure ahead of the wave $p_0\approx 1$ atm. The diffusion loss of electrons from the laser beam during the time τ_L is quite insignificant in this case: in addition, $\omega_0\gg\nu_c^2$, and therefore all the terms in (5) are proportional to the pressure and there is no dependence on p_0. The results shown in Fig. 1 are valid not only under the conditions of the experiments of Refs. 3–5, but in all cases when the following two conditions are satisfied: $\omega^2\gg\nu_c^2$ and $(D_e\tilde{\nu}^{-1})^{1/2}\ll R$. The abscissas represent the CO_2-laser radiation intensity, but if $\omega^2\gg\nu_c^2$ the same results, as seen from (5), are obtained for a neodymium laser having an intensity 100 times larger. The noticeable heating of the atoms and ions begins only after the degree of ionization of the plasma z exceeds a value $\sim 10^{-2}$ (Refs. 6,7), and it was therefore assumed in the calculation of the avalanche that $T=0$.

The threshold for the development of a fast ionization wave is a laser intensity q_{thr} such that during the time of the laser pulse the plasma front is only insignificantly detached from the initiating laser spark. The electron avalanche develops slowly, and the duration of the laser pulse suffices only to ionize a thin gas layer adjacent to the spark. Let n_1 be the number of electrons per unit volume of the plasma behind the ionization-wave front. At $q=q_{thr}$, the condition $n\sim n_1$ is satisfied by the instant $t=\tau_L$ therefore Eq. (3) can be represented in the form

$$\tau=\tau_L[\ln(n_1/\dot{n}_{ph}\tau)]^{-1}. \tag{6}$$

Eliminating $\tilde{\nu}$ from (5) with the aid of (4) and substituting (6), we obtain a simple expression for the threshold:

$$q_{thr}=\frac{mc}{4\pi e^2}(\omega^2+\nu_c^2)\times\left[\nu_c^{-1}\left(\bar{I}+\frac{3}{2}T_{e0}\right)\left(\tau_L^{-1}\ln\frac{n_1}{\dot{n}_{ph}\tau}+\frac{D_e}{L^2}\right)+\frac{3m}{M}(T_{e0}-T)\right]. \tag{7}$$

This expression is similar to the expression for the breakdown threshold q^* of the pre-ionization gas. The difference is that the density of the priming electrons n_0 is replaced by the product $\dot{n}_{ph}\tau$, so that the difference between the threshold q^* and q_{thr} is determined by the value of the logarithm. In the case of breakdown of fresh gas without pre-ionization, however, the breakdown threshold q^* is determined not by the losses but precisely by the difficulty of producing the initial density n_0 (via multiphoton photoionization of the atoms or the tunnel effect), and an expression of the type (7) for q^* is not valid at all. This manifests itself particularly clearly in experiment with a CO_2 laser: pure argon does not break down even at $q>1$ GW/cm² (see, e.g., Ref. 10, p. 65); in the case of preionization at a level $n_0\sim 10^{11}$ cm³, the breakdown threshold is $q^*\approx 200$ MW/cm² (Ref. 10), and the threshold for the onset of fast ionization wave is $q_{thr}\approx 30$ MW/cm² (Ref. 3).

The value of $\dot{n}_{ph}$ is calculated in the next section (see Eq. (10)). It is determined by the parameters of the laser spark. In the experiments of Refs. 1–5 the spark parameters were not measured, but for all reasonable parameter combinations the logarithm in (7) is of the order of 10. For long laser pulses ($\tau_L\gtrsim 10^{-4}$ sec), the exact value of the logarithm is of no importance at all, since the first term is much smaller than the other two. Such a relation between the parameters was satisfied in the experiment of Ref. 2. The results of calculation by means of Eq. (7) are in good agreement with the results of this experiment (Fig. 2). In experiments with a CO_2 laser[3–5] the radiation pulse was much shorter, therefore, the threshold q_{thr} was determined mainly by the first term of (7). Not knowing the parameters of the "priming" plasma, only the logarithm can be estimated, and this determines the threshold with accuracy to within a factor of the order of unity. The time between the flash of the priming plasma and the termination of the laser pulse depended on the energy input[3–5,14]

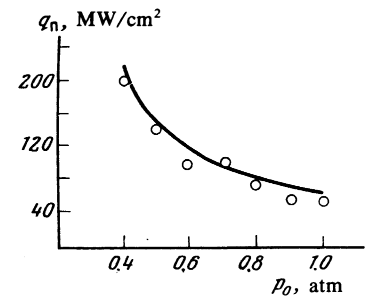

FIG. 2. Dependence of the threshold of the onset of the fast ionization wave on the initial xenon pressure. Calculation by formula (10) for the experimental conditions of Ref. 2, and the experimental results (points). Neodymium laser.

at the threshold of the onset of the fast ionization wave, this time is $\sim$40 nsec. The threshold amounted to (30–40) MW/cm^2 for argon and xenon at $p_0 = 1$ atm (Ref. 3). According to (7), $q_{\rm thr} \approx 30$ MW/cm^2 (xenon) and $q_{\rm thr} \approx 60$ MW/cm^2 (argon).

In Ref. 15 were investigated optical-detonation and radiation waves in xenon. The intensity of the radiation of the neodymium laser did not exceed 400 MW/cm^2. The threshold of the onset of the fast ionization wave (7), calculated for the conditions of this experiment, was $400/p_0$ MW/cm^2. In the experiment $p_0 \leqslant 1$ atm, so that the fast ionization wave could not propagate. Earlier experiments by the same authors were carried out in air.[16] The radius of the beam was smaller and the radiation intensity reached 2.5 GW/cm^2. At $p_0 = 0.5$ atm, a transition from the optical-detonation regime of the plasma-front propagation to the superdetonation regime was observed. The authors interpret the latter as radiative, although the value $u = 115$ km/sec obtained for the velocity should have been observed (according to their theory of the radiation wave) for intensities five times larger. In our opinion, the strong growth of $u(q)$ is evidence of a transition to a fast ionization wave. It is precisely near the threshold that it is characterized by a strong dependence $u/u_{\rm thr} = (q/q_{\rm thr})^a$ with $a \gg 1$. For a radiation wave the condition would be $a < 1$.

VELOCITY OF NONSTATIONARY MOTION OF THE PLASMA FRONT

If the laser-radiation intensity exceeds the threshold for the onset of the fast ionization wave, the plasma front begins to move along the laser beam, becoming detached from the bare laser spark. The front velocity at the start of the path can differ considerably from the velocity of the steady-state motion of the wave. In the former case the gas ahead of the front of the wave is pre-ionized by the laser-spark radiation, and in the latter it is pre-ionized by radiation of the plasma layer produced in the front. The spectral composition and intensity of the ionizing radiation are substantially different in the two cases, since the plasma temperatures, the degrees of ionization, and the sizes of the radiating volumes are all different. The steady-state velocity of the fast ionization wave was calculated in Refs. 6 and 7. The present section is devoted to calculation of velocity during the initial stage of the front motion. In the investigation of the plasma-front motion regimes, measures were taken to produce a plane and homogeneous front, in particular, by using laser beams of cross section. This circumstance causes only a negligible part of the electrons to be fused from the beam during the time of the laser pulse, and according to (4) we have $\tau \approx \tilde{\nu}^{-1}$. The solution (3), starting with the instant $t \approx \tilde{\nu}^{-1}$, can be represented in the form

$$n(x,\, t) = \dot{n}_{\rm ph}(x)\, \tilde{\nu}^{-1} e^{\tilde{\nu} t}. \tag{8}$$

To calculate the wave velocity, we select inside the front a plane on which the number of free electrons per unit volume reaches a given value $n_a \lesssim n_1$. The width of the front is many times smaller than the path traversed by it, therefore the front velocity u, even for a nonstationary wave, is practically independent of the plane chosen inside the front to determine its coordinate. The condition $n(x,t) = n_a$ and expression (8) make it poissible to calculate the plasma-front law of motion $x = x(n_a, t)$. Differentiating with respect to time, we obtain

$$u = \dot{x}(n_a,\, t) = \left(\frac{dn}{dt}\right)_a \left(\frac{dn}{dx}\right)_a^{-1} = \tilde{\nu}\dot{n}_{\rm ph}(x) \left(\frac{d\dot{n}_{\rm ph}}{dx}\right)^{-1}, \tag{9}$$

where the index a means that the derivative is calculated at $n(x,t) = n_a$.

The photoionization velocity $\dot{n}_{\rm ph}(x)$ in the gas surrounding the laser spark can be determined from Eq. (1). We write down the radiation intensity of the plasma in the plane-layer approximation and use the known expression for the total absorption coefficient in a multiply ionized plasma.[11] for typical values of the parameters characterizing the laser spark, its optical thickness in the hard part of the spectrum ($h\nu > I_i$) is small. Expanding the exponential and correcting the absorption coefficient to allow for the stimulated emission, we obtain

$$\dot{n}_{\rm ph}(x) = 1.4\cdot 10^8 R^* N_Z (Z+1)^2 T \exp\left\{-\frac{I_{Z+1}}{T}\right\} \times \int_{I_i}^{\infty} \varkappa_\nu^{\rm ph,0} \exp\{-\varkappa_\nu^{\rm ph,0} x\} \frac{d\nu}{\nu} \quad [\mathrm{cm}^{-3}\mathrm{sec}^{-1}]. \tag{10}$$

Here N_Z is the number of Z-ions per cm^3, and I_{Z+1} is their ionization potential. All the quantities in front of the integral pertain to the spark, R^* is expressed in centimeters, T in electron volts, and $\varkappa$ in cm^{-1}.

Differentiating (10), we can represent the velocity in the form

$$u(x) = \tilde{\nu}(q) N_0^{-1} \psi(N_0 x), \tag{11}$$

where $\psi \equiv \varphi_1/\varphi_2$, while

$$\varphi_{1,2}(N_0 x) \equiv \int_{I_i}^{\infty} \sigma_{\rm ph}^{1,2}(\nu) \exp[-N_0 x \sigma_{\rm ph}(\nu)] \frac{d\nu}{\nu}, \tag{12}$$

N_0 is the number of atoms in a unit volume of gas ahead of the front, and $\sigma_{\rm ph}(\nu)$ is the atoms photoionization cross section. The minus sign that specifies the direction of the veloc-

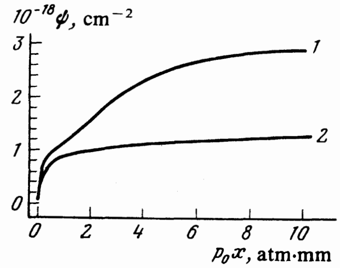

FIG. 3. The function $\psi(N_0x)$. Argon (1) and xenon (2). $N_0 = 2.7\cdot10^{19}p_0$ cm^{-3}. The gas pressure p_0 is in atmospheres.

ity is left out. For multielectron atoms, the cross section σ_{ph} is a rather complicated nonmonotonic function of the frequency, making it difficult to simplify expression (11) further. The integrals (12) depend only on the parameter N_0x. The function $\psi(N_0x)$ can be calculated for each gas by integrating the experimentally measured cross sections. By way of example, Fig. 3 shows the results obtained for argon (curve 1) and xenon (curve 2). We used the cross-section tables given in Refs. 17–19.

Equation (11) and Figs. 1 and 3 enable us to determine the velocity of the fast ionization wave on the initial section of the path. The frequency $\tilde{\nu}$ of the inelastic collisions of the electron is proportional to N_0, so that the ratio $\tilde{\nu}N_0$ is independent of pressure—it determines only the dependence of the velocity on the laser-radiation intensity.

The three experimentally observed regularities are explained by the form of the function $\psi(N_0x)$. First, the wave velocity increases rapidly with increasing distance from the priming plasma.[5] This growth of $u(t)$ is observed even against the background of the decreasing intensity $q(t)$ at $q\approx50$ MW/cm^2 (argon, 1 atm). The strong dependence of ψ on N_0x near the priming plasma causes the detachment of the wave front from the initial spark to have the character of an abrupt ejection—the acceleration du/dx is large. Second, experiment has revealed a decrease of the wave velocity with decreasing pressure p_0 (Ref. 5), which also follows from the form of the function $\psi(N_0)$. Third, at a considerably lower pressure ($p_0\lesssim0.1$ atm) no fast ionization wave is produced, since the velocity of its front near the spark in this regime would be lower than the velocity of the front in the optical-detonation or in the radiative regimes. It must be noted here that the decrease of the front velocity with decreasing p_0 is a distinguishing feature of the unsteady motion of the fast ionization wave. All the steady-state regimes are characterized by a growth of the front velocity with decreasing p_0.

With increasing distance of the front of the fast ionization wave from the bare laser spark, a smaller role is assumed by the photoionization produced by the priming plasma, and at a distance $x\gg(N_0\sigma_{ph})^{-1}$ from the spark, the main source of the ionizing radiation becomes the plasma layer adjacent to the front. The plasma in the beam is produced as the fast ionization wave propagates, and its radiation leads to a further advance of the front, namely, the wave reaches a stationary propagation regime (if the laser pulse duration is long enough).

DEPENDENCE OF THE FRONT VELOCITY AND OF THE PLASMA TEMPERATURE ON THE LASER-BEAM RADIUS

We continue the investigation of the steady-state motion of a fast ionization wave, initiated in Refs. 6 and 7. Let a plane wave front propagate uniformly opposite to the laser radiation. A superdetonation regime exists, therefore the expansion of the plasma begins at a distance $\gtrsim R$ behind the front. The plasma temperature prior to expansion is T_1 and the density is N_0.

In the investigation of the regimes of the motion of the plasma front, the functions $u(q)$ and $T_1(q)$ are usually approximated by power laws in the form

$$u\sim q^a,\quad T_1\sim q^b.$$

This description is quite illustrative when the regimes are compared, but it is not rigorous, since the exponents themselves are heat functions of the laser-radiation intensity. Constant values of the exponents a and b correspond to an approximation in which the effective adiabatic exponent is independent of the temperature and of the degree of ionization of the plasma. Such an approximation become valid at a temperature comparable with the ionization potential of a hydrogenlike ion, when $\gamma\to5/3$.

A fast ionization wave is characterized by exponents $a>1$ and $b<0$. For the remaining plasma-front propagation regimes we have $a<1$ and $b>0$. Another distinguishing feature of a fast ionization wave is the dependence of the front velocity and of the plasma temperature behind the front on the laser-beam radius. By way of example, Fig. 4 shows results obtained for xenon in a CO_2 laser beam. The intensity of the laser radiation was assumed constant, $q(t)=$ const. The results were obtained by the same method as used in Refs. 6 and 7, namely numerical integration of the equations that describe the ionization kinetics, the laser-radiation absorption, and the heating of the electron and ion subsystems; the wave velocity was obtained as an eigenvalue of the problem.

The calculations performed have revealed a number of new regularities that characterize the steady-state motion of a fast ionization wave. If the laser beam radius is small, the wave velocity increases with increasing radius. As R increases, the growth of the velocity slows down gradually and stops eventually. The growth of the velocity stops at small R if N_0 is large, or at large R if N_0 is small. At $N_0\geqslant10^{20}$ cm^{-3}, the wave velocity and the plasma temperature behind the front are independent of the laser-beam radius if $R>10^{-2}$ cm. The temperature $T_1(q,N_0,R)$ is a decreasing function of the three indicated arguments. Lowering of the plasma temperature with increasing laser-radiation intensity is a distinguishing feature of a fast ionization wave; the cause of this dependence was established in Ref. 7. We shall show that the remaining regularities that characterize plasma-front propagation regime also follow from simple reasoning.

Once the motion of the fast ionization wave steadied, the source of the ionizing radiation is a plasma layer of thickness $\sim R$ adjacent to the front. The intensity of the ionizing radiation is described by an expression of the form

$$J_\nu(x)=J_{\nu p}(T_1)\,[1-\exp(-\varkappa_{\nu 1}'R)]\,\exp(-N_0\sigma_{ph\,\nu}x)$$

$\varkappa_{\nu1}'=\varkappa_\nu'(T_1)$ is the spectral absorption coefficient corrected

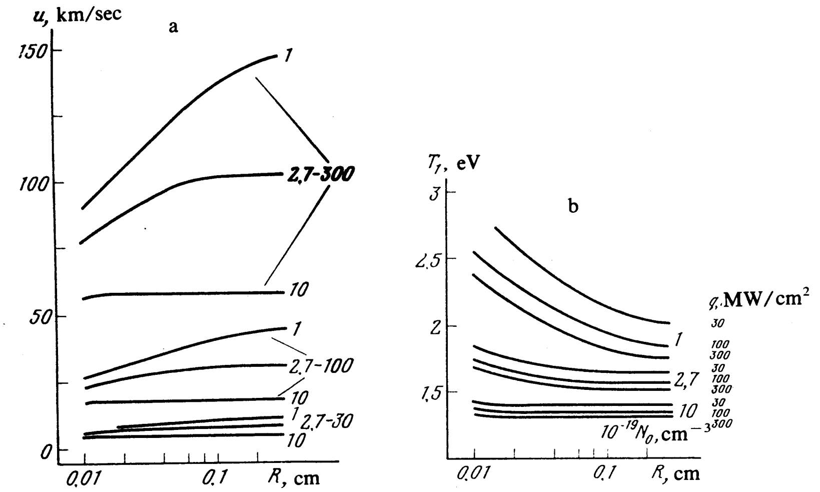

FIG. 4. Dependence of the velocity of a fast ionization wave (a) and of the plasma temperature behind the front (b) on the laser-beam radius. Xenon. CO_2 laser. The numbers represent $10^{-19}N_0(\text{cm}^{-3})$ and q (MW/cm²).

for the stimulated emission). For further estimates it is necessary to separate a spectral interval in which the photons "trigger" the electron avalanche. The function $\sigma_{ph}(\nu)$ is not monotonic and varies quite considerably in magnitude, so that the intensity $J_\nu(x,T_1,R,N_0)$ also depends on the frequency in a rather complicated manner. The principal role in the "cutting out" of the indicated interval played by two circumstances: 1) at $h\nu > 3T_1$, the intensity $J_{\nu eq} \sim \exp(-h\nu/T_1)$ of the equilibrium radiation decreases exponentially with increasing $h\nu$; in xenon $T_1 < 3$ eV, and the decrease of $J_{\nu eq}$ is very strongly dependent on the photoionization threshold itself; 2) the cross section $\sigma_{ph}(\nu)$ has a local minimum in the interval (40–60) eV, where $\sigma_{ph}(\nu) \equiv \sigma_{min} = (1.5–2)\cdot 10^{-18}$ cm².

As a result of these circumstances, the "triggering" of the electron avalanche is effected by photons of energy $h\nu > 40$ eV (xenon). In the hard part of the spectrum ($h\nu > 160$ eV), the cross section σ_{ph} is even smaller, but there the intensity $J_{\nu eq}$ is already extremely small. The range of the photons with $h\nu \approx 40$ eV in the cold gas ahead of the front is $l_{40} \sim (N_0\sigma_{min})^{-1}$ and amounts to ~ 0.2 mm at $p_0 = 1$ atm. A numerical calculation shows that the electron avalanche is triggered at a distance $\lambda_1 \approx (3–5)l_{40}$ ahead of the wave front.

The optical thickness of the radiating plasma layer in this part of the spectrum is $\Lambda_{40} \sim N_0\sigma_{min}AR$. Here A is the fraction of the atoms from among the N_0 atoms and ions. So long as the condition $\Lambda_{40} \ll 1$ is satisfied, the radiation intensity is proportional to Λ_{40} and increases in proportion to R. The condition Λ (40 eV) = 1 makes it possible to determine that critical value of the beam radius $R_\Lambda \sim (N_0\sigma_{min}A)^{-1}$ above which the wave velocity and the plasma temperature T_1 are independent of R. The fraction of the atoms is determined from the Saha equation. In the calculation performed (Fig. 4), the temperature $1.3 < T_1 < 2.7$ eV and the fraction of atoms in the plasma is $0.1 < A < 0.9$.

We consider two laser beams with different radii (R' and R''), but with equal radiation intensity. If $R' < R'' < R_\Lambda$, we have $\dot n'_{ph}(x) < \dot n''_{ph}(x)$ and $\lambda'_1 < \lambda''_1$, i.e., in the laser beam with the larger radius the electron avalanche will be triggered at a larger distance from the wave front. However, $\tilde\nu' = \nu''$ at equal laser radiation intensity; therefore both avalanches will evolve in the same time after the triggering, and during this time the plasma fronts will cover different distances ($\lambda'_1 < \lambda''_1$), i.e., $u' < u''$. This means that the velocity of the fast ionization wave should increase with increasing R, until its radius reaches a value of the order of R_Λ, after which the radius has practically no effect on $\dot n_{ph}$ and the velocity becomes independent of the radius.

From the energy conservation law

$$N_0u[{}^3/_2T_1(1+z_1)+Q(z_1)]=q$$

($z_1 \equiv z(T_1)$ is the equilibrium ionization degree) and from the condition $u' < u''$ at $q' = q''$ is follows that $T'_1 > T''_1$, meaning that the plasma temperature behind the fast ionization wave should decrease with increasing radius ($R < R_\Lambda$). Here $Q(z)$ is the energy lost to ionization.

With decreasing gas pressure and laser-beam radius, the plasma temperature behind the wave front increases, and the fraction of the atoms decrease; at $T_1 > 2.7$ eV the radiating plasma layer becomes optically so thin that no fast ionization wave is produced.

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