DYNAMICS OF THE PROPAGATION OF A FAST WAVE OF IONIZATION OF A GAS IN A LASER BEAM

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Several regimes of propagation of optical breakdown in gases exposed to laser radiation have been studied thus far both experimentally and theoretically. These regimes correspond to different ranges of intensities of the incident radiation and pressures of the undisturbed gas medium and differ from one another by the mechanisms of transfer of the plasma front along the light channel [1-3]. In recent years researchers have been interested in the fast ionization wave (FIW) regime [3-8], in which the motion of the breakdown wave is determined by the thermal emission of the plasma, corresponding to the development of an electron avalanche in the field of the laser radiation in front of the light absorption front [3]. This regime is characterized by a very sharp dependence of the velocity of the ionization front u on the intensity of the incident radiation q ($u \sim q^a$, $a > 1$), which distinguishes it from the well-known regimes for which $a < 1$.

For all well-known regimes of propagation of an optical breakdown wave the stationary stage has been studied in greatest detail theoretically. In so doing, one specific mechanism for transfer of the plasma front, determining the basic characteristics of the process, stands out. However, the stationary analysis may be inapplicable in studies with short pulses or in cases when the intensity of the incident radiation is close to the threshold intensity for a definite breakdown propagation regime. In such situations the competition between different mechanisms of transfer of the ionization front is important (especially at the starting stage of the motion) and the characteristic time at which the stationary regime is reached may be comparable to the duration of the laser pulse. It is obvious that in this case the process is not stationary, and it must be studied on the basis of nonstationary gas dynamics, which takes into account the entire complex of physical factors affecting the flow of the process [9].

This work is devoted to the study of the nonstationary interaction of laser radiation with gases. For simplicity we shall study hydrogen, for which it is possible to follow all characteristic features of the process and for which the most complete data on the atomic and kinetic coefficients are available in the literature. The range of radiation intensities and initial gas densities adopted in the calculations is chosen so that it would be possible to follow the character and the regularities of the process accompanying the transition from the photodetonation regime of propagation of a plasma front to the FIW regime.

Formulation of the Problem. Basic Equations. Unlike the stationary phase of the process, in which stationary relations between the intensity of laser radiation and the characteristics of the absorption wave at an arbitrary moment in time are studied, we shall study these regularities at the stage of formation of the breakdown wave from the moment at which the plasma "seed" forms [10, 11]. As a rule, breakdown is initiated either by a preliminary short pulse with a radiation intensity exceeding the threshold for breakdown of the gas or by interaction of the laser beam with a barrier [4, 6, 12]. In both cases the starting "seed" is quite transparent to the incident radiation, and the formation of the breakdown wave, as a comparison of calculations with experimental photoregistrograms show, has qualitatively the same character. We shall study a "seed" in the form of a local plasma formation in the gas medium with a characteristic size $L < \ell_\omega$ (ℓ_ω is the mean-free path length of laser quanta). The main laser pulse is switched on at the time $t = 0$ and it starts the evolution of the plasma region. In the process, as experiment shows, the breakdown wave propagates only in the channel of the light beam and the geometry of the process is close to a planar one-dimensional case [6].

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The determining factors in the formation of the breakdown wave and its structure and dynamic characteristics are the absorption of radiation, electron–ion energy exchange, and ionization and radiation transfer. Because of the difference in the relaxational properties of the plasma components at the wave front the discontinuity between the electron and ion temperatures as well as the deviation of the degree of ionization and the intensity of the plasma radiation from equilibrium values could be significant. As a rule, such processes are studied on the basis of two-temperature, one-dimensional gas dynamics, taking into account all physical factors enumerated above.

Based on the foregoing, the system of equations is written in the form

$$\frac{\partial}{\partial t}\left(\frac{1}{\rho}\right)-\frac{\partial v}{\partial m}=0,\quad dm=\rho dx,\quad \frac{\partial v}{\partial t}+\frac{\partial (p_i+p_e)}{\partial m}=0,$$
$$v=\frac{dx}{dt},\quad \frac{\partial E_i}{\partial t}+p_i\frac{\partial v}{\partial m}=\frac{1}{\rho}Q_{ei},$$
$$\frac{\partial E_e}{\partial t}+p_e\frac{\partial v}{\partial m}=-\frac{1}{\rho}Q_{ei}-\frac{\partial W}{\partial m}+\frac{\partial q}{\partial m}+q_T, \qquad (1)$$
$$\rho\frac{\partial q}{\partial m}=\varkappa_\omega^e q,\quad W=-\chi\rho\frac{\partial T_e}{\partial m},$$
$$q_T=-2\pi\frac{1}{\rho}\int_0^{\mu_m}d\mu\int_0^\infty \varkappa_\nu^*(I_{\nu p}^*-I_\nu)\,d\nu,\quad \mu\rho\frac{dI_\nu}{dm}=\varkappa_\nu^*(I_{\nu\,\mathrm{eq}}^*-I_\nu)$$

(for convenience in performing the numerical calculations, Lagrangian variables are employed). Here ρ is the density; v is the mass velocity; $Q_{ei} = (3/2)(m_e/m_i)n_e\nu_{ef}(T_e - T_i)$ is the collisional energy transfer from electrons to ions; T_e and T_i are the electron and ion temperatures; and, ν_{ef} is the effective frequency of elastic collisions between electrons and heavy particles.

Since in the interaction process the plasma in the zone of radiation absorption consists of partially ionized gas, collisions of electrons with ions and neutral particles contribute to ν_{ef}, and $\nu_{ef} = \nu_{ei} + \nu_{en}$, where $\nu_{ei} = 2(2\pi/m_e)^{1/2}e^4\Lambda n_iT_e^{-3/2}$, Λ is the Coulomb logarithm, and n_i is the ion density. The expression from [13] can be used for ν_{en}: $\nu_{en} = 2.2\cdot 10^{-7}N_a$ (N_a is the density of neutral particles).

The heat flux W is determined by the electronic thermal conductivity $\chi(T_e, \rho)$, which is calculated taking into account also the interaction of electrons with the neutral components of the plasma and can be taken in the form [10]

$$\chi^{-1}=\chi_e^{-1}+\chi_a^{-1},$$
$$\chi_e=0.75\frac{\gamma_0 n_e}{\Lambda e^4 Z^2 n_i}(2\pi m_e)^{-1/2}T_e^{5/2},\quad \chi_a=0.28\frac{n_e}{N_a}\left(\frac{8T_e}{\pi m_e}\right)^{1/2}\frac{1}{\sigma},$$

where Z is the ion charge, $\gamma_0 = 3.16$ (for hydrogen), σ is the total electron-neutral collisions cross section, and n_e is the electron density.

The internal energy and pressure of the electronic and ionic components of the plasma equal, respectively,

$$E_i=\frac{R}{\gamma-1}T_i,\quad p_i=R\rho T_i,\quad E_e=\frac{R}{\gamma-1}zT_e+Q(z),\quad p_e=R\rho zT_e. \qquad (2)$$

Here Q(z) is the specific ionization energy, while $z = n_e/N$ is the degree of ionization determined from the kinetic equation

$$dz/dt=f(z,T_e,\rho) \qquad (3)$$

(the explicit form of the right side is written out below).

We shall now consider in greater detail how the thermal radiation of the plasma is taken into account in the model under study. In the photodetonation regime the front of the wave of absorption of laser radiation is transferred hydrodynamically. The plasma radiation has virtually no effect on this process and can be taken into account in the equation of conservation of energy as a loss [11]. In the case of FIW the thermal radiation, as is obvious from estimates given in [2], also does not participate in the heating of the gas and its role reduces exclusively to creating the seed ionization in front of the boundary of the plasma region in the zone of the light beam, which leads later to development of an electron avalanche and transfer of the zone of light absorption in the field of the incident wave. Thus the transfer of thermal radiation can be studied only along the laser light channel with radius R_1, assuming for simplicity that the intensity I_ν is uniform over the entire transverse section of the beam. This means that I_ν is a function of only two variables: the coordinate x and the angle ϑ between the direction of propagation of the radiation and the channel axis ($\mu = \cos\vartheta$), and in addition the region of variation of ϑ, as follows from the geometry under study, is also a function of x ($0 \le \vartheta \le \vartheta_{max}(x)$).

The spectral coefficient of absorption of thermal radiation $\varkappa_\nu$ consists of the coefficient of bound-free absorption $\varkappa_\nu^{ph}$, responsible for the photoionization of atoms by hard quanta, and the spectral coefficient of absorption of radiation owing to elastic collisions of electrons with ions and atoms $\varkappa_\nu^e$ ($\varkappa_\nu = \varkappa_\nu^{ph} + \varkappa_\nu^e$). The expression for $\varkappa_\nu^{ph}$ has the form† [14] $\varkappa_\nu^{ph} = \frac{32\pi^2 Z^2 e^6}{3^{3/2} h^4 c \nu^3} N_a I_H$, $h\nu > I_H$, and the coefficient $\varkappa_\nu^e$, corrected for induced emission, can be taken in the form

$$\varkappa_\nu^e = \frac{4\pi\sigma_\nu}{cn_\nu}\frac{T_e}{h\nu}\left(1-\exp\left(-\frac{h\nu}{T_e}\right)\right), \tag{4}$$

where I_H is the ionization potential of hydrogen; σ_ν is the conductivity of the plasma; n_ν is the refractive index; $\sigma_\nu = \frac{e^2 n_e}{m_e}\frac{\nu_{ef}}{\nu^2+\nu_{ef}^2}$; $n_\nu^2 = 1 - \frac{\omega_{pe}^2}{\nu^2+\nu_{ef}^2}$; $\omega_{pe}^2 = \frac{4\pi n_e e^2}{m_e}$.

In the case of laser radiation $\hbar\omega \ll I_H (\omega = 2\pi\nu)$ and its absorption is described only by the expression (4), which for $\hbar\omega < T_e$ transforms into $\varkappa_\omega^e = \frac{\omega_{pe}^2}{c}\frac{\nu_{ef}}{\omega^2+\nu_{ef}^2}\left(1-\frac{\omega_{pe}^2}{\omega^2+\nu_{ef}^2}\right)^{-1/2}$.

The electron density, which appears in the expression for the coefficients of thermal conductivity, the absorption of radiation and electron–ion exchange in Eqs. (1), is determined from (3). The electron density changes in two independent channels: inelastic interaction of electrons with atoms and photoionization processes. The rate of photoionization equals $\frac{dn_e^{ph}}{dt} = 2\pi\int_0^{\mu_m} d\mu \int_{I_H/h}^{\infty} \frac{\varkappa_\nu^{ph} I_\nu}{h\nu}\,d\nu$. In the process, electrons whose energy can be substantially higher than the ionization potential I_H are created. As a result of inelastic collisions such photoelectrons subsequently lead to the appearance of additional electrons, and in this manner the initial thermal quantum generates more than one electron. It is shown in [2] that for laser radiation fluxes sufficient to realize the FIW regime it may be assumed that each act of photoionization leads effectively to the appearance of two thermal electrons, so that in (3) a factor of 2 must be included in the terms responsible for the photoionization processes. The rate of impact ionization is taken in the standard form [14] $\alpha = Cv_e(I_H + 2T_e)\exp(-I_H/T_e)$, $v_e = (8T_e/\pi m_e)^{1/2}$. Recombination constants are calculated from the ionization rates in terms of the equilibrium constants $b_{ph} = a_{ph}/K_{eq}$, $\beta = \alpha/K_{eq}$, $a_{ph} = (1/N)dn_e^{ph}/dt$.

Thus (3) can be written in the form

$$\frac{dz}{dt} = \frac{dz_{ph}}{dt} + \frac{dz_e}{dt} = 2a_{ph} - 2b_{ph}Nz^2 + \alpha Nz - \beta N^2 z^3, \tag{5}$$

where z_e and z_{ph} correspond to electrons produced by impact ionization and photoionization.

Since the medium in the zone of radiation absorption is not in a state of thermodynamic equilibrium, the deviation of the charged-particle densities from their equilibrium values must be taken into account in the coefficients in the radiation transfer equation. Constructing the kinetic equation for the radiation intensity function I_ν taking into account the above-discussed mechanisms of radiation absorption in a nonequilibrium medium, we obtain the corrected radiation absorption coefficient $\varkappa_\nu^*$ and the "equilibrium" radiation intensity $I_{\nu eq}^*$:

†Here and below we assume that the ionization of an atom occurs from the ground state. This approximation is valid in the region of low temperature in front of the absorption wave. In the region of high temperatures the role of photoionization is insignificant. The validity of this approximation for the shock-wave mechanism is demonstrated in [15].

$$\varkappa_\nu^* = \varkappa_\nu^e + \left[\frac{1 - \delta \exp\left(-\frac{h\nu}{T_e}\right)}{1 - \exp\left(-\frac{h\nu}{T_e}\right)} \right] \varkappa_\nu^{ph}, \quad I_{\nu eq}^* = \frac{\varkappa_\nu^e + \delta\varkappa_\nu^{ph}}{\varkappa_\nu^*} I_{\nu eq}, \quad \delta = \frac{z(1 - z_{eq})}{z_{eq}(1 - z)}.$$

Here z_{eq} is the equilibrium degree of ionization, determined from Saha's equation, and $I_{\nu eq}$ is the spectral intensity of the equilibrium radiation. It is easy to verify that for $z = z_{eq}$ the equation of radiation transfer has the usual form [14].

Results. In performing the calculations the initial value of the degree of ionization of the undisturbed gas was set equal to $z_0 \approx 0$ and the temperature $T_0 = 0.2$ eV. For simplicity the gas was assumed to be atomic. Taking into account the kinetics of dissociation complicates the calculations somewhat, but does not significantly change the results.

The calculations performed revealed definite regularities in the emergence of the ionization wave into the stationary state. For characteristic parameters of the "seed" (L ~ 50-100 μm and $T_e = T_i = 10$ eV) it is quite transparent to the incident radiation and absorbs only part of the laser energy. The plasma region at first expands owing to heat conduction by electrons; in the process, because of the sharp pressure gradient, a shock wave begins to form in the gas adjacent to the "seed." At the same time, under the action of the radiation from the hot region a precursor of the ionization front, in which the electron temperature breaks away from the ion temperature, forms in the cold gas. For the foregoing parameters of the "seed" the electron temperature increases up to some constant value T_e^*, determined by the intensity of the incident radiation q, within approximately 0.3 nsec (for a neodymium glass laser with λ = 1.06 μm) and this increase is virtually independent of the gas density and q. In the process, the spatial variation of T_e and T_i, starting at some distance from the plasma front, is so insignificant that their profile may be regarded as practically constant over the entire volume of the precursor (Fig. 1, q = 30 GW/cm²). The ionization is also insignificant and does not exceed $z \sim 10^{-4}$-10^{-3}.

As the temperature increases, increasingly more of the incident radiation is transmitted through the "seed," but owing to heating and ionization of the adjacent mass of gas the optical thickness of the region occupied by the ionized gas increases and the zone of absorption of radiation becomes gradually localized at the shock wave. In the process, the change in the characteristics of the hot region of the plasma has virtually no effect on the parameters of the precursor, which are determined by the insignificant characteristic absorption of laser radiation in this region. This structure also remains unchanged later, when the ionization wave detaches from the "seed," and can be explained based on the following considerations. In the region of the undisturbed gas in front of the plasma front at a distance of the order of the mean-free path length of thermal quanta, which make the greatest contribution to photoionization, the laser energy absorbed by the electrons goes into heating the particles and impact ionization of the gas. For $z \ll 1$ $\nu_{ef} \sim$ const, therefore $\varkappa_\omega \sim z$ and $Q_{ei} \sim z$, and in addition $dz/dt \approx \alpha N z$. Taking into account (2) and (5), the equations describing the energy balance for this region can be written as

$$\partial T_e/\partial t = Aq - [B(T_e - T_i) + Q\alpha N], \quad \partial T_i/\partial t = Bz(T_e - T_i). \tag{6}$$

It is obvious from here that for $z \ll 1$ $\partial T_i/\partial t \simeq 0$ and T_i remains constant and approximately equal to T_0. In addition, T_e also reaches some value T_e^*, depending on q, for which $\partial T_e/\partial t \sim 0$. Near the plasma front where z increases sharply the expression (6) is no longer valid; the contribution of other energy redistribution mechanisms becomes important, T_e and T_i increase, and as a result there forms a radiation absorption zone with a structure characteristic for the corresponding regime of propagation of an optical breakdown wave. It is obvious that the precursor forms in front of any emitting plasma front absorbing laser radiation, so that there arises the question: what ultimately determines the character of the propagation of the ionization wave? This question can be answered qualitatively as follows. The thermal radiation creates seed electrons in front of the plasma front in a channel of length $l \sim 1/\varkappa^{ph}$. The characteristic time for the development of an avalanche in the gas $\tau_1 \sim 1/N\alpha$, i.e., the boundary of the ionized gas moves toward the beam with a velocity $u \sim N\alpha/\varkappa^{ph}$. If the velocity of the radiation absorption zone D, associated, say, with the transport of hydrodynamic disturbance, exceeds u, then there is not enough time for an avalanche to develop in the undisturbed gas and the motion of the plasma boundary occurs by a hydrodynamic mechanism, i.e., the regime is a photodetonation regime. If, however, there is enough time for an avalanche to form ($u \geq D$), then the motion of the plasma boundary is determined by the velocity u, and in the process the gas behind the ionization wave remains

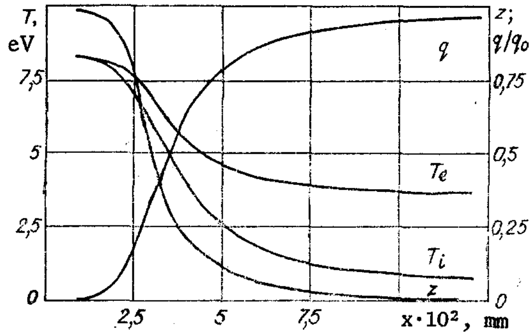

Fig. 1

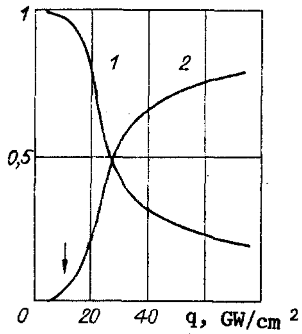

Fig. 2

undisturbed, which corresponds to the FIW regime. Thus a natural criterion is obtained for the transition from the photodetonation regime to FIW: $N\alpha/\varkappa^{ph} \geq D$. The arguments presented here are consistent with [3]: the transfer of the front occurs by the mechanism for which, under specific conditions, the velocity of the plasma boundary is highest.

Following the criterion obtained above, we can attempt to evaluate the threshold intensity of the incident radiation q_{th}, after which the FIW regime appears. For this we return to the formulas (6). The laser energy absorbed in the precursor is expended primarily on ionization and heating of the gas. Since in this region T_i and T_e do not change significantly, and $T_e = T_e^*(q)$, the role of these mechanisms in energy redistribution can be estimated. Figure 2 shows the dependence of the relative fraction of the energy expanded on heating the ions (curve 1) and ionization of the gas (curve 2) in the precursor on q. It is obvious that for $q \simeq q_{th}$ (the threshold value q_{th} is marked with an arrow) $Q_{ei} \gg Qdz/dt$, so that $Aq \simeq B(T_e - T_i)$. Since $A = \varkappa_\omega \rho^{-1} \sim \rho\omega^{-2}$ and $B \sim \rho$, $q\omega^{-2} \sim T_e$. The rate of ionization $\alpha \sim T_e^{1/2}\exp(-I_H/T_e)$, while $D = [2(\gamma^2 - 1)q/\rho]^{1/3}$. Substituting all this into the expression for the criterion and taking the logarithm, we obtain $q_{th} \simeq A^*\omega^2(\ln q_{th} + 2\ln\rho - 6\ln\omega + C)^{-1}$, where A^* and C are constants. This expression, obtained from qualitative considerations, correctly describes the dependence of the threshold intensity on the laser frequency and the weak dependence, found computationally, of the threshold on the starting gas density.

Finally, we would like to point out one other fact. The computed dependence of the velocity of the ionization front on the incident radiation for some value $q^* > q_{th}$ has a discontinuity: $u \sim q^{2-4}$ for $q_{th} < q < q^*$ and $u \sim q$ for $q > q^*$. This is also confirmed by physical experiments [7]. The change in the character of FIW on both sides of q^* can be most clearly followed for the dependence of the plasma temperature behind the wave front on q [2]. The value of q^*, as calculations showed, is virtually independent of the gas density and, for example, equals approximately 25-30 GW/cm² for hydrogen in the case of neodymium laser radiation. This behavior of the curve u(q) can also be qualitatively explained, based on the fact that one of the mechanisms for dissipation of the energy absorbed in the precursor plays the dominant role. Turning once again to Fig. 2, it is obvious that the energy flux from electrons to ions and the ionization losses Qdz/dt are also equal to one another when $q \simeq 26$ GW/cm². For $q > q^*$, $Qdz/dt > Q_{ei}$ and we can set $q \simeq Q\alpha N/A$. From here $u \sim N\alpha/\varkappa^{ph} \sim q$. In the case $q < q^*$, $q \sim T_e$, arguing analogously we obtain $u \simeq Kq^{1/2}\exp(-I_H/T_e)$, which differs from the expression $u \sim q^a$, which approximates the experimental and computed data. However, taking the logarithm of both these expressions, differentiating with respect to to ln q, and equating the right sides we obtain (the derivatives of the neglected constant factors vanish)

$$a = 1/2 + I_H/T_e. \tag{7}$$

It follows from both experiments and calculations [2] that the index a assumes the value ~4 at the threshold of the regime and decreases as the intensity of the incident radiation q increases. The same behavior also follows from the formula (7). As the calculations performed show, the electron temperature T_e in the precursor equals ~3 eV for q = 20 GW/cm². This gives an index a ~ 5. As q increases in the interval $q_{th} < q < q^*$ the electron temperature in the precursor increases, which causes a to decrease.

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FEATURES OF IONIZATION AND EMISSION BEHIND STRONG SHOCK WAVES IN AIR

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In 1960-1970 numerous experiments were carried out for studying emission processes in air (see, e.g., [1]) on the basis of shock tubes. Studies embraced a wide spectral range (λ = 200-6000 nm), and equilibrium values of temperature 2000-14,000 K with gas pressure $p_0 > 13$ Pa. On the basis of experiments the role of different emission processes was established and values of parameters required for quantum mechanics calculations for emission were specified. Results of detailed calculations of emission equilibrium for air are presented in tables in [1, 2].

By comparing calculated values for spectral intensity of emission for air from these experimental works it is possible to draw the following conclusions. The main mass of experimental results obtained in shock tubes relate to values $T \le 10^4$ K and $\rho \ge 10^{-3}\,\rho_0$ (ρ_0 is density under normal conditions, ρ is gas density). There is good conformity between results of experiments and calculations. A more complex situation is observed in analyzing data found at higher temperatures and low values of ρ. A considerable proportion of experiments with $T > 10^4$ K were carried out in shock tubes with observation of emission in the region behind the reflected shock wave (SW) with relatively high gas density and pressure. With low pressures and $T > 10^4$ K the study of emission capacity was carried in a few works (e.g., [3, 4]). In them experimental conditions were not analyzed for explaining the presence behind the shock wave of local thermodynamic equilibrium (LTE). It was assumed that in the region behind the SW, corresponding to the output of the recorded signal of the emission receiver at a quasi-steady level, LTE exists. Values of spectral emission intensity obtained under these conditions (λ = 500, 6100 nm) with $v_s > 11$ km/sec are markedly below the corresponding calculated values (if the calculation is carried out using [1, 2]).

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