Superdetonation motion of plasma front towards a powerful laser beam

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We continue the investigation of a fast gas-ionization wave propagating along a laser beam of prebreakdown intensity [Sov. Phys. JETP **52**, 1083 (1980)]. An explanation is presented of a number of experimental results [V. V. Korobkin *et al*., Sov. Phys. JETP **26**, 79 (1968); I. Z. Nemtsev and B. F. Mul'chenko, Sov. J. Plasma Phys. **3**, 649 (1977); V. A. Boĭko *et al*., Sov. J. Quant. Electron. **8**, 134 (1978); V. D. Zvorykin, Paper at All-Union Conf. on Interaction of Optical Radiation with Matter, Leningrad, 1981].

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INTRODUCTION

The motion of a plasma front towards a laser beam was studied in many experiments and investigated theoretically (see, e.g., Refs. 1–6, the literature cited in Ref. 1, and Raĭzer's monograph[7]). Some of the 1980 experimental results,[2-4] however, could not be explained on the basis of the prevailing ideas concerning the conditions of the motion of the plasma front and of the associated laser-radiation absorption wave. Nemtsev and Mul'chenko[3] have shown that the regime in their experiment cannot be identified with any of the known regimes.[7] This new regime was named fast ionization wave (FIW). Also observed was an unexpectedly strong (much stronger than linear) dependence of the front velocity u on the laser-emission intensity q_m (Refs. 2–4), whereas all the calculations predicted a $u(q_m)$ weaker than linear.[1)] The experimental results in Refs. 2–4 pointed to the existence of at least one regime with properties different from those of all those hitherto investigated.

This conclusion prompted one of us to investigate further the possible plasma-front transport mechanisms.[1] Calculations have shown that at a laser-emission intensity exceeding a certain threshold $\hat{q}_m$ it is possible to observe a regime that differs greatly in its properties from all those studied before. The threshold of the regime, calculated for the conditions of the experiments of Nemtsev and Mul'chenko,[3] is close to observed one, and the new regime[1] is therefore identified with the FIW,[3] and will so designated hereafter. The FIW is in fact characterized by a strong $u(q_m)$ dependence. If the dependence of the front velocity and of the plasma temperature T^* behind the front on the laser emission intensity is approximated by powerlaw functions

$$u-\hat{u}\sim(q_m-\hat{q}_m)^a,\quad T^*-\hat{T}^*\sim(q_m-\hat{q}_m)^b, \tag{1}$$

then the FIW is characterized by exponents $a>1$ and $b<0$ (Ref. 1), thereby distinguishing it from the other regimes, for which these inequalities are reversed.[7] The threshold values are marked by carets. The condition $a>1$ corresponds to a strong (stronger than linear) $u(q_m)$ dependence, while the condition $b<0$ corresponds to a decrease of the plasma temperature behind the FIW front with increasing laser intensity, a most unusual behavior.

The main task in the initial stage of the FIW investigation[1] was to identify the physical mechanism capable of ensuring a strong $u(q_m)$ dependence. The calculations were made for monatomic hydrogen, since the cross sections for the elementary processes in this gas have been most thoroughly investigated. The experiments, however, were not performed with hydrogen, so that the theoretical and experimental results on the FIW agree only so long as the condition $a>1$ is satisfied, and the realization of the proposed front-transport mechanism remains an open question.

The basic investigations aimed at observing FIW[4-6] were made for argon and xenon, using a short (nanosecond) laser pulse.[2)] The first purpose of the present paper is the calculation of the FIW as applied to the conditions of these experiments. Agreement between the calculations and the experimental results can be regarded as weightly evidence favoring the proposed plasma-front transport mechanism.[1] In addition, we continue the investigation of the FIW; the problem is reformulated to take into account the finite duration of the laser pulse, and an expression is obtained for the threshold of the regime.

PLASMA-FRONT TRANSPORT

The thermal radiation of the plasma propagates in the cold gas ahead of the absorption front, in the direction counter to the laser emission. The hard part of the thermal radiation ($h\nu>I$) ionizes certain atoms, and if the laser intensity is high enough, the photoelectrons detached from the atoms within the boundaries of the light beam initiate an electron avalanche. Here I is the ionization potential of the atom. With increasing distance from the luminous front of the discharge, the intensity of the ionizing radiation is decreased by the absorption. To obtain a near-unity probability of initiating an electron avalanche at a certain distance ahead of the front, the thermal radiation must ensure here a definite physically small but finite rate of appearance of priming electrons, $\dot{n}_{ph}^{min}$. A reasonable order of magnitude of $\dot{n}_{ph}^{min}$ can be determined from the experimental conditions. In the present case, for example, one can choose a rate at which one photoelectron is produced in a volume of 10^{-6} cm^3 within 0.1 nsec, i.e., $\dot{n}_{ph}^{min}=10^{-16}$ cm$^{-3}\cdot$sec^{-1}. The distance λ^* between the discharge front and the parallel thin cold-gas layer, where the photoionization

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rate is $\dot{n}_{ph}^{min}$, can be calculated from the known dimensions of the discharge and the spatial distributions of the thermodynamic variables.

The motion of the plasma front in the FIW regime is the result of weak (priming) photoionization of the gas at a distance λ^* ahead of the front. This is in essence the only role of the ionizing radiation, inasmuch as the subsequent gas ionization is by avalanche in the laser-emission field. As the fresh gas layers become ionized, the absorption zone moves in a direction counter to the laser emission, and its glow generates priming electrons in new gas layers a distance λ^* ahead of the luminous front.[3)] The avalanche evolution time τ_{av} is determined by the laser-emission intensity, and the charge propagation velocity is close to $\lambda^*(T^*)/\tau_{av}(q_m)$.

Calculations show that only during the initial growth stage of the degree z of the ionization does the photoionization rate $\dot{z}_{ph}$ exceed the rate $\dot{z}_e$ of the impact ionization, since $\dot{z}_e \sim z$. The main growth of the degree of ionization takes place under the condition $\dot{z}_e \gg \dot{z}_{ph}$.

The difference between the FIW and the radiation wave[7] is that the FIW only "triggers" the avalanche, and the energy needed to ionize the gas comes from the laser emission. The radiation wave heats and ionizes the gas until effective absorption of the laser radiation sets in, i.e., $T_0 \sim 2$ eV and $z \sim z_{eq}(T_0)$ (Ref. 7, p. 222), where z_{eq} is the equilibrium degree of ionization. This difference between the degrees of ionization of the gas by the radiation, inherent in the model of the phenomenon,[1,7] leads to fundamental differences in the laws governing the discharge propagation. Thus, for example, in the case of the radiation wave an increase of the plasma temperature with increasing q_m is assumed to be obvious (Ref. 7, p. 219), while for the FIW the calculations yield a decrease of $T^*(q_m)$.

The ionization wave produced in argon and xenon when a CO_2 laser beam acts on a graphite partition was investigated experimentally in Refs. 4–6. The radiation pulse was a spike of 120 nsec duration at half maximum with a flat peak (~60 nsec) and a long tail containing half the energy. The recorded plots show that when the radiation intensity $q_0(t)$ approaches the maximum value q_m the ionization wave propagation becomes stationary. The front is almost plane, inasmuch as at a beam diameter ~2 cm the wave moves several millimeters away from the partition during the time of the spike. The experimental conditions thus conform to the approximations made by us in Ref. 1, when steady-state motion of a plane ionization wave was considered.

FORMULATION OF PROBLEM

We discuss in this section the changes that must be made in the problem as formulated in Ref. 1 to take into account the actual experimental conditions of Ref. 4–6.

Allowance must be made first for the finite duration τ_L of the laser pulse. In the case of a sufficiently long laser pulse, the integration of the system of equations describing the FIW started from the coordinate $x = \lambda^*$ determined by the discharge dimensions and by the spatial distribution of the thermodynamic variables. A degree of ionization $z(\lambda^*) = 0$ was specified at the point $x = \lambda^*$, and the derivative $\dot{z}(\lambda^*)$ was $\dot{z}_{ph}^{min} = \dot{n}_{ph}^{min}$ in accord with the definition of λ^*; here N_0 is the gas density ahead of the wave front. In the superdetonation regime of the plasma-front propagation, the expansion of the gas begins behind the zone where the laser emission is absorbed. The velocity of the FIW is an eigenvalue of the problem. The method of determining $u(q_m)$ is described in detail in Ref. 1. If the inequality $\lambda^* > u\tau_L$ is satisfied at a certain laser intensity, the laser-pulse time will not be long enough for the absorption-wave front to negotiate the distance λ^*, i.e., at a given q_m the avalanche cannot develop fully during the time τ_L if the initial photoionization rate is $\dot{z}_{ph}^{min}$. For arbitrary τ_L, the initial coordinate must obviously be chosen to be

$$\lambda = \begin{cases} \lambda^*, & \text{if} \quad \tau_L > \lambda^*/u \\ \tau_L u, & \text{if} \quad \tau_L < \lambda^*/u \end{cases},$$

and it is necessary to specify $z(\lambda) = 0$ at this point. We have in this case for the derivative $\dot{z}(\lambda) > \dot{z}(\lambda^*) \equiv \dot{z}_{ph}^{min}$, since $\lambda < \lambda^*$.

The time of establishment of the FIW regime in the experiment differs little from the time τ_1 of formation of a luminous plasma layer at the surface of the partition. The laser intensity approaches q_m within this time. A plot of $\tau_1(q_m)$ is shown in Fig. 3 of Ref. 6. The FIW regime terminates at the instant $\tau_2 \approx 140$ nsec, when the laser intensity drops to $q_2 \approx 0.7 q_m$. We assume hereafter that the laser pulses that maintain the steady motion of the FIW are rectangular with intensity $q_0(t) = q_m$ and duration

$$\tau_L(q_m) = \tau_2 - \tau_1(q_m). \tag{2}$$

The width $H(t)$ of the luminous plasma layer at the target surface increases with time, but the formation of the priming electrons is influenced primarily by the initial plasma luminosity, when the layer width is a fraction of a millimeter.

Comparing the thermal velocity of the electron with the velocity of its oscillations in the light-wave field, it can be shown that the electron energy distribution function $f(\varepsilon)$ begins to deviate substantially from Maxwellian at $q_m \sim 1$ GW/cm^2 (CO_2 laser) or 100 GW/cm^2 (neodymium laser). The FIW was observed in experiment at much smaller q_m, so that the frequencies of the ionization and of the excitation of the atoms by electron impact, as well as the frequencies of the elastic collisions of the electron with the atoms and with the ions, were calculated using a Maxwellian distribution function. The frequencies of the inelastic processes are the most sensitive to the form of $f(\varepsilon)$; we shall therefore calculate the error due to replacement of the true distribution function by a Maxwellian for the excitation rate constant α_e^* (Ref. 8). We compare the $\alpha_e^*(E/\omega)$ dependence calculated by Phelps[9] for the true distribution function with the analogous dependence that can be easily obtained from Eq. (20) of Ref. 1, using explit expressions for the frequencies and assuming $f(\varepsilon)$ to be Maxwellian. Here ω and E are the frequency of the laser emission and the field strength in the light wave. The upper curve of Fig. 1 was obtained for the following conditions: the excited atoms are ionized immediately, the

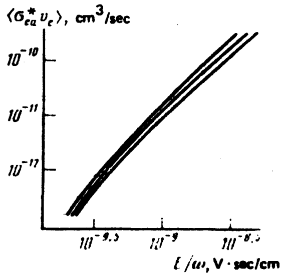

FIG. 1. Rate of excitation of atoms by electron impact vs the field intensity in the light wave. Argon, for details see the text.

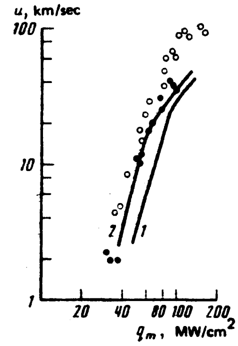

FIG. 2. FIW velocity in argon (1, ○) and xenon (2, ●). Initial pressure 1 atm. CO_2 laser. Experimental results of Ref. 4.

probability of direction ionization is considerably lower than the probability of ionization via excited level, and the electron energy is equal to the potential of the first excited level. The lower curve was obtained with account taken of direct ionization of the atom by electron impact and of ionization via excited levels, with the electron energy equal to the ionization potential. The middle curve was taken from Phelps's paper. The relative error due to the use of a Maxwellian distribution function does not exceed 20% in terms of the field strength.

The rate constants of the ionization and excitation of atoms by electron impact were calculated in the usual manner (see Ref. 8) from the experimental data given in Refs. 7, 8, and 10. The effective frequency ν_{ea} of elastic collisions of the electron with the atoms was calculated by integrating the experimentally measured cross sections (see Ref. 11). The results of the integration, with accuracy no less than 5%, are approximated by the expressions

$$\nu_{ea}=10^{-7}n_a(1.3T_e-0.7-0.1T_e^2)\ \text{sec}^{-1},$$
$$\nu_{ea}=10^{-7}n_a(4.8T_e-2.65-0.69T_e^2)\ \text{sec}^{-1}$$

for argon and xenon respectively ($1 \leqslant T_e \leqslant 6$ eV). The cross section $\sigma_{ph}(h\nu)$ for the ionization of the atoms were taken, in the numerical solution of the problem, from the detailed tables of the experimental results.[12,13]

DISCUSSION OF THE RESULTS

The $u(q_m)$ relation obtained by numerically integrating the system of equations of Ref. 1 is compared with the experimental results in Fig. 2. The agreement for xenon is almost complete, while for argon the results differ by a factor of almost 1.5 in the radiation intensity. The cause of this discrepancy is not clear at present. The main result, however, is the equality of the exponent a in Eq. (1) for the observed and calculated relations. This agreement holds for both argon and xenon, even if account is taken of a deviation from the power law, i.e., if $a=a(q_m)$. Thus, the discussed plasma-front transport mechanism yields the correct character of the $u(q_m)$ dependence. All other transport mechanisms yield a much smaller slope of the curve.

Great interest attaches also to Fig. 3; the plasma temperature behind the FIW front decreases with increasing q_m, i.e., with increasing energy input to the laser pulses. In other absorption-wave propagation regimes the plasma temperature rises with increasing laser intensity. Let us clarify the cause of so unusual a form of the function $T^*(q_m)$. The degree of ionization in the avalanche increases exponentially:

$$z(t)=z_0 e^{\tilde{\nu}t}$$

right up to near the equilibrium value $z_{eq}(T_e)$. Here $\tilde{\nu}(q_m)$ is the frequency of electron production, and $z_0(T^*)$ is the priming degree of ionization produced by the plasma radiation upon photoionization of the atoms from from a distance $x \approx \lambda$.

Consider rectangular laser pulses of duration τ_1 and τ_2 and respective intensity q_1 and q_2. Let $q_1 < q_2$, then $\tau_1(q_1) < \tau_2(q_2)$ according to (2). The electron temperature in the avalanche is also higher in the second case and $\tilde{\nu}_2 > \tilde{\nu}_1$. In both cases, one value $z_{abs} \gtrsim 0.01$, corresponding to the start of the effective absorption of the laser emission, i.e.,

$$z_{01}e^{(\tilde{\nu}_1\tau_1)}=z_{02}e^{(\tilde{\nu}_2\tau_2)}, \tag{3}$$

but the product $\tilde{\nu}_1\tau_1 < \tilde{\nu}_2\tau_2$ and to satisfy Eq. (3) we need $z_{01}(T_1^*) > z_{02}(T_2^*)$, i.e., $T_1^* > T_2^*$ at $q_1 < q_2$. It turns out thus that a lower temperature behind the plasma front is necessary to transport the front in the FIW regime at a higher laser intensity. The true picture is somewhat more complicated than this simple scheme, for λ increases whem q_m increases, so that a decrease in the priming photoionization is due in part to the increased distance from the luminous layer; the main decrease of z_0 is nevertheless connected with the lowering of the temperature.

The transport of the plasma front in the FIW regime is not connected with motion of the matter. The energy conservation law

$$q_m=N_0u[{}^3/_2T^*(1+z^*)+Iz^*]$$

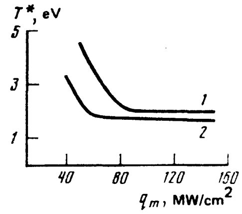

FIG. 3. Maximum electron temperature behind FIW front. Curve 1—argon, 2—xenon. Initial pressure 1 atm. CO_2 laser.

explains the strong $u(q_m)$ dependence. With increasing energy input into the laser pulses, T^* decreases and the degree of ionization $z^*(T^*)$ of the plasma behind the front decreases. To conserve the energy in this case, the function $u(q_m)$ must be stronger than linear.

We now describe briefly the structures of the FIW in argon and xenon. The electron temperature ahead of the absorption-wave front, in the electron avalanche development zone (i.e., $0 < x \leqslant \lambda$, $z_{abs} \geqslant z \geqslant 0$) is constant and is determined by the laser-emission intensity, more accurately by the equilibrium between the absorption and the losses. When the ionization approaches equilibrium, the inelastic losses decrease and T_e begins to increase. The profile of the electron temperature has the usual maximum in the zone where the laser radiation is absorbed. The energy is transferred from the electrons to the atoms and to the ions slowly because of the large difference between the masses of the particles. In the absorption region, the temperature of the heavy particles is still much lower than 1 eV. The greater part of the laser radiation is absorbed in a plasma layer whose thickness is several microns ($N_0 = 2.7\times10^{19}$ cm^{-3}). The electron and ion temperatures are equalized in argon and in xenon over a length of the order of several dozen microns. In hydrogen, the relaxation took place (mainly) within the limits of the absorption zone. The degree of ionization of the plasma at the start of the ionization zone is much lower than the equilibrium value, and approaches the latter towards the end of the zone.

THRESHOLD OF THE REGIME

When the laser intensity is reduced, the electron temperature in the avalanche $T_{e0}(q_m)$ is lowered, the frequency of the inelastic collisions falls off exponentially, and the plasma-front transport velocity is decreased. It is easy to determine the laser-emission intensity $\hat{q}_m$ at which the energy absorbed by the electron gas suffices only to compensate for the elastic losses:

$$\hat{q}_m = 80\Omega^2 T_{e0} A^{-1}\ \text{MW/cm}^2,\quad \omega^2 \gg \nu_c^2. \tag{4}$$

Here ν_c is the frequency of the elastic collisions of the electrons with the atoms and the ions, Ω is the ratio of the laser frequency to that of the CO_2 laser, A is the atomic weight of the gas, and $T_{e0}\approx 1$ eV. At $q_m < \hat{q}_m$ there is no transport of the plasma front, i.e., $\hat{q}_m$ is the strict (unattainable) threshold of the regime.

Motion of the plasma front in the FIW regime was observed experimentally for xenon in a neodymium laser beam at $q_m \approx 100$ MW/cm^2 (Ref. 3). The $\hat{q}_m$ threshold (4) under these conditions is 60 MW/cm^2. For argon in a CO_2 laser beam, the threshold is $\hat{q}_m = 2$ MW/cm^2; no measurements were made at $q_m > 20$ MW/cm^2, inasmuch as at low laser intensity no priming plasma layer (whose emission would "divert" the FIW from the surface) managed to be formed at the graphite surface during the time of the laser pulse. In the experiment with neodymium laser[3] the laser pulse duration was of the order of microseconds, as against several dozen nanoseconds in the experiments with the CO_2 laser, so that in the former case it was possible to come much closer to the regime threshold. In both cases, motion of the plasma front was observed at a laser intensity that was low compared with the gas optical-breakdown threshold q^*.

The FIW threshold is quite high for hydrogen in a neodymium-laser beam, $\hat{q}_m \approx 8$ GW/cm^2. For many gases at atomspheric pressure this intensity is sufficient to produce optical breakdown. In pure hydrogen, however, under conditions when the priming electrons can appear only as a result of 12-photon ionization of the atoms by the laser radiation, the breakdown threshold is $q^* \sim 500$ GW/cm^2. This value can be easily obtained on the basis of the results of Bebb and Gold.[14]

In conclusion, we dwell briefly on certain results obtained in Refs. 2, 3, and 5. A jumplike plasma-front motion was observed in the case of sharp focusing of the neodymium laser beam in air.[2] The threshold $\hat{q}_m$ increased with decreasing gas pressure.[3] The front velocity varied in a complicated manner (but not greatly) when the pressure was lowered from 1 to 0.001 atm.[5] In the former case the exponent $a \approx 1.5$ in Eq. (1) points to FIW, but within the framework of the existing model one can claculate only the steady-state motion of the front, while the stationary problem (which is not planar because of the sharp focusing) calls for another formulation. In the second and third cases, the wave becomes semitransparent with decreasing pressure. The fraction of the laser radiation absorbed in the wave front could not be measured, and this fraction is itself dependent on q_m. In addition, in the case of a semitransparent absorption wave the method used to initiate the wave (auxiliary flash[3] or a partition[5]) influences the laws governing the front transport during the time of the entire laser pulse. The appropriate mathematical problem is nonstationary, and is difficult to solve in a formulation that takes the experimental conditions into account.

At the same time, at an initial pressure $p_0 \gtrsim 1$ atm and higher the wave is opague to the laser radiation and the formulation of the problem is relatively simple. Experiments at increased gas pressures are of interest. They should reflect the dependences of the regime threshold, of the velocity, and of the temperature on the gas pressure, the beam radius, and the laser-pulse duration. In this case the laws governing the FIW differ from those for other regimes.

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[1)] The waveform of the laser pulse prior to its absorption in the medium is $q_0(t)$. The maximum instantaneous emission intensity in the pulse is $q_m = \max q_0(t)$.

[2)] The authors of Refs. 4—6 identified the observed wave with the radiation wave, the most likely from among those known at that time. They have indicated, however, that the results do not agree with the $u(q_m)$ relation obtained in Ref. 7 for radiation waves.

3) The charge propagation is assumed here stationary, and the laser pulse long enough.

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