

SCREENING OF A SURFACE VAPORIZING UNDER  
LASER BOMBARDMENT, AT TEMPERATURE AND  
IONIZATION DISEQUILIBRIUM

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The kinetics of the process through which vapor released from the surface of a solid is heated by laser radiation are calculated. The data indicate the existence of temperature and ionization disequilibrium (for Al at incident flux densities  $q$  over  $50 \text{ MW/cm}^2$  in times of the order of  $1 \mu\text{sec}$  or less). Treatment of this disequilibrium lowers the "critical" flux density  $q_*$  corresponding to the onset of screening within a specified time  $t$ , or within the time  $t_*$  required for screening to develop at a specified flux density  $q$ , appreciably, as compared to the values obtained in the equilibrium theory of the process. Some of the other physical processes that might have some effect on "flare" behavior are also discussed.

The effect of laser radiation on the surface of a solid can be brought about under two very distinct sets of conditions:

1) Vapor possessing a temperature  $T$  close to the temperature  $T_V$  of the phase transition, vapor weakly ionized and practically transparent to the incident radiation (evaporation wave with no "screening," no evaporating surface [1-4]);

2) Vapors are heated to  $T \gg T_V$ , are strongly ionized (at least singly ionized), and absorb a substantial portion of the incident radiation, only a small part of which manages to reach the evaporating surface (the surface is thereby "screened" [4-7]).

Since the distribution of the parameters in the layer of vapor, the patterns of the variations with time  $t$ , and the characteristic values of the parameters per se (the temperature  $T$ , the pressure  $p$ , the density  $\rho$ , and the rate of outflow  $u$ , etc.) differ substantially depending on the conditions under which the effect is brought about, it would be of interest to find out more about how a transition from one set of conditions to the other is effected.

The phenomenon of breakdown in cold gases through the action of a bombarding laser beam is quite familiar (a review of the literature may be found in [8]). One of the present authors [9, 10] developed a theory of the equilibrium flare. The flare differs from breakdown first in the fact that the vapor is already appreciably ionized because of the sufficiently high vapor temperatures. For instance, at  $T_V$  values corresponding to  $p$  of the order of  $10^2$  to  $10^3$  bar, a typical value of the degree of ionization  $\alpha = 10^{-4}$  (of the total number of particles present). The number  $\nu_{ee}$  of collisions between electrons and the number  $\nu_{aa}$  of collisions between ions, and also the number  $\nu_{ia}$  of ion-atom collisions, is therefore such that we can speak of the existence of an electron temperature  $T_e$ , an atom temperature  $T_a$ , and an ion temperature  $T_i$ .

For example, when aluminum is targeted by laser radiation of flux density  $q = 100 \text{ MW/cm}^2$ , the data for the vapor are  $T = 4500^\circ\text{K}$ ,  $\rho = 10^{-2} \text{ g/cm}^3$ ,  $\alpha = 0.45 \cdot 10^{-3}$ ,  $\tau_{ee} = 1.5 \cdot 10^{-13} \text{ sec}$ , which is much shorter than the characteristic time  $\tau$  in which the radiation exerts its effect, or the flare development time  $t_*$ . The transient period of the atom temperature at those same parameters  $\tau_{aa} = 1.5 \cdot 10^{-12} \text{ sec}$ , and we obtain  $1.5 \cdot 10^{-12} \text{ sec}$  for ion-atom collisions. Consequently, the electron temperature  $T_e$ , the atom temperature

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$T_e$ , and the ion temperature  $T_i$  are all established within a time  $\approx 10^{-12}$  sec, with  $T_i$  practically equal to  $T_e$  (so that we can assume in fact  $T_e = T_i$  in the remainder of the discussion).

The quantities  $T_e$  and  $T_i$  are set equal by two distinct mechanisms. The first, which plays an important role at high degrees of ionization, is that of collisions with ions. The characteristic electron-ion relaxation time  $\tau_{ei} = 5 \cdot 10^{-9}$  sec. The second mechanism, which dominates in the initial stage of the process, involves encounters with atoms, and the electron-atom relaxation time  $\tau_{ea} = 5 \cdot 10^{-10}$  sec. At the same time, the typical duration  $\tau$  of the laser action in the Q-spoiling mode ranges from  $3 \cdot 10^{-9}$  to  $3 \cdot 10^{-7}$  sec [6, 7]. The characteristic duration  $\tau$  of a single spike in the free-running mode is  $10^{-6}$  sec [1, 3]. Consequently, a considerable portion of the energy absorbed by the electrons is used up not only in heating a comparatively small number of electrons, but also in raising the temperature of the entire volume of gas as a whole (raising the ion temperature and the atom temperature). We can assume, in a first crude approximation,  $T_e = T_i = T$ , and the degree of ionization  $\alpha$  is then determined by the equilibrium temperature  $T$  of the material according to Saha's equation. That assumption was in fact entertained in [9, 10]. There the flare is seen to develop as follows. Since the variable  $\alpha$  (when  $\alpha \ll 1$ ) depends exponentially on the temperature, with a large exponent (inasmuch as the first ionization potential  $I_1 \gg T$ ), a slight rise in the temperature is sufficient for a marked increase in  $\alpha$ . For example, when  $I_1/T = 10$ , it is sufficient to raise  $T$  by 10% for  $\alpha$  to about triple in value, so that the absorptivity  $\kappa$ , and with it the rate of increase in  $T$ ,  $\alpha$ ,  $\kappa$ , etc., will increase correspondingly (the warm-up process is "self-accelerating" in nature, with a sudden absorption increase resulting).

In contrast to breakdown, this flaring phenomenon takes place in moving material. Expansion of vapor into a rarefied medium or into a vacuum is accompanied by adiabatic cooling of the vapor, which can even exhibit heating because of absorption of the radiation. As the thickness of the layer of vapor increases and as the pressure gradient declines, the role played by adiabatic cooling gradually tapers off, while the rate of energy release on account of absorption of radiation remains unchanged (at constant flux density of the oncoming radiation, and if we ignore transfer of heat to the bulk of the body by ordinary heat conduction). Consequently, the role played by those two factors is commensurate at any given moment, and the warm-up process described above then sets in. The flare thus shows a critical character under conditions where the material is in complete equilibrium.

Yu. V. Afanas'ev and O. N. Krokhin [11] proceeded, in estimating the time of onset of substantial warm-up of the vapor, from the view that this warm-up process would have to set in with intense absorption of radiation, which would occur when the optical thickness was of the order of unity, or more precisely  $\tau = 0.28$ , i.e., the same as in the approximate theory of self-consistent conditions [5].

The theory behind the appearance of the flare [9, 10] is based on a local criterion of the equality of the amount of energy released through absorption of radiation and through cooling due to disintegration. This criterion can also be reduced to  $\tau = \tau_*$  when the flux is constant, but  $\tau_*$  depends on the heat-transfer properties of the materials such as the heat of evaporation  $Q$ , the enthalpy of the vapor in the phase transition  $h_v$ , and the reflection coefficient  $k_r$ . At typical values of those parameters, we find that the flare occurs in vapors that are practically transparent, i.e.,  $\tau_* \ll 1$  (e.g., aluminum when the optical thickness  $\tau_*$  is of the order of 0.01).

The computed values of the flux density  $q_*$  at which the flare occurs in hot vapor are not only much less than the corresponding values needed for breakdown in cold gases, but are also well below the estimated values in [11].

However, as we can see by comparing the experimental results cited by A. I. Petrukhin [12] and the data predicted by V. I. Bergel'son on the basis of theory [9, 10] and run on a computer for the real shape of the pulse occurring in [12], the screening phenomenon sets in at lower  $q_*$  value than anticipated from the theory. We therefore have to turn again to an analysis [13] of the initial assumptions in [9, 10], and most of all the assumption that the vapor is close to a state of temperature equilibrium and ionization equilibrium.

According to the estimates in [14], under conditions where flux densities  $q \leq 10^3$  MW/cm<sup>2</sup>, as was the case in [12], the divergence between the electron temperature and the atom temperature remained small:  $\Delta = T_e - T_i \ll T_e \approx T$ . But the criterion for the weak effect of that divergence must not be the inequality shown, because of the clearly defined exponential dependence  $\alpha(T)$ , but must be some stronger inequality  $\Delta(I_1/T) \ll T$ . Moreover, the increase in  $T_e$  and in  $\alpha$ , and accordingly also in  $\kappa$ , brings about a still greater increase in  $\Delta T$ .

Note that that type of effect had already been pointed out [15] in estimates made of the onset of screening of a radiating strong shock wave by a gas ahead of the shock front, which is also, as would appear from the appropriate experiments, observed earlier than predicted by the pure equilibrium theory [16].

1. The system of equations describing heating of electrons by radiation with account taken of losses in the work of adiabatic expansion and energy transfer to atoms and ions (in collisions), as well as heating of ions and atoms through such energy transfer, appears in the form

$$\begin{aligned} (c_v)^e \frac{dT_e}{dt} + J + p_e \frac{1}{v} \frac{dv}{dt} &= f - (Q_{ea} + Q_{ei}) \\ [(c_v)^i + (c_v)^a] \frac{dT_i}{dt} + (p_a + p_i) \frac{1}{v} \frac{dv}{dt} &= (Q_{ea} + Q_{ei}) \end{aligned} \quad (1.1)$$

Here  $(c_v)^e$ ,  $(c_v)^a$ ,  $(c_v)^i$  are the respective specific heats of the electrons, atoms, and ions, and  $p_e$ ,  $p_a$ , and  $p_i$  are their respective partial pressures. Since estimates of the relaxation times of the degree of ionization to its equilibrium value are indicative of how validly they can be compared to the characteristic screening times  $t_*$ , equations (1.1) are supplemented here by the kinetics equation of the ionization process

$$\frac{dn_e}{dt} = \alpha_p n_a n_e - \beta_p n_i n_e^2 \quad (1.2)$$

where  $n_e$ ,  $n_a$ ,  $n_i$  are the respective number densities of electrons, atoms, and ions per  $\text{cm}^3$ ;  $\alpha_p$  and  $\beta_p$  are the respective constants of the ionization rate and recombination rate.

We now introduce a few simplifying assumptions, none of which are fundamental in nature, pertinent to the terms appearing in Eqs. (1.1) and (1.2). The first terms in equations (1.1) express energy losses in raising the temperature of the electrons, ions, and atoms. The term  $J$  appearing in Eq. (1.1) takes into account energy losses by electrons in ionization. Electrons escaping an atom have a negligible thermal velocity, and some additional energy is still required in order to heat them up to the electron temperature. In considering the process by which screening develops, we limit the discussion to temperatures  $\approx 0.3$  to 2 eV, and take into account the fact that

$$J = \left( I_1 + \frac{3}{2} T_e \right) \frac{dn_e}{dt} \quad (1.3)$$

where  $I_1$  is the first ionization potential.

The terms  $p_e dv/vdt$  and  $(p_i + p_a) dv/vdt$  describe cooling of electrons, ions, and atoms through adiabatic separation of the particles.

Let us now examine in greater detail the terms  $Q_{ei}$  and  $Q_{ea}$  describing exchange of energy between electrons and ions, and between electrons and atoms. The relaxation time for attaining a Maxwellian distribution between the charged particles can be found in [16]. In the particular case of electrons and singly charged ions, we have

$$\tau_{ei} = \frac{3.5 \cdot 10^8 A T_e^{7/2}}{n_i Z^2 \ln \Lambda} \quad (1.4)$$

where  $A$  is the atomic weight of the ions,  $Z$  is the charge on the ions, and for the parameters assigned in the problems under consideration  $\ln \Lambda \approx 6$ ,  $n_i$  is the number of ions per  $\text{cm}^3$ ,  $T_e$  is the electron temperature in eV, and  $\tau_{ei}$  is in sec.

For the relaxation time of the temperature in electron-atom collisions, we have

$$\tau_{ea} = \frac{m_a}{n_a \langle v_e \rangle \sigma 2m_e} \quad (1.5)$$

Here  $\langle v_e \rangle$  is the mean velocity of the electrons;  $\sigma$  is the gas-kinetic cross section for elastic collisions;  $\sigma = \pi r^2$ , where  $r$  is assigned equal to the dimension of the atom in the Thomas-Fermi model in the case of heavy elements, and equal to the radius of the atom in the Bohr model in the case of light elements;  $r = na_0 \sqrt{I_H/I_1}$ ,  $I_H = 13.6$  eV,  $n$  is the principal quantum number of an electron in the outer shell in the ground state of the atom;  $a_0$  is the Bohr radius;  $m_e$  is the mass of the electron;  $m_i$  is the mass of an ion. Two exchange mechanisms are present when temperatures are not sufficiently high, and the total relaxation time will exhibit the form

$$1/\tau = 1/\tau_{ei} + 1/\tau_{ea}$$

For the flux of energy from electrons to ions and atoms per  $\text{cm}^3$ , we use the expression

$$Q_{ei} + Q_{ea} = 3/2 n_e (T_e - T_i) 1/\tau \quad (1.6)$$

The term in equations (1.1) describing absorption of radiation by electrons assumes, in the assumption of almost complete transparent vapor, the form

$$f = (1 + k_r) q (\kappa_i + \kappa_n) \quad (1.7)$$

$\kappa_i$  is the mass absorption coefficient for free-free electron transitions in the field of the ion and bound-free transitions from highly excited states of the atom, and  $\kappa_n$  is the bremsstrahlung absorption coefficient in collisions with a neutral atom. Bremsstrahlung absorption was taken into account, when calculating the value of  $\kappa_i$ , under the assumption that the Gaunt factor was unity, while the photoeffect from highly excited states obeyed the hydrogen-like Unsold-Kramers approximation with summation over states replaced by integration [16]. In that approach, the value of  $\kappa_i$  is  $\exp(\epsilon/T)$  times in excess of the coefficient for bremsstrahlung absorption on ions. We note, for example, that absorption takes place, in the case of Al and a ruby laser, starting from the excited level with the principal quantum number  $n = 4$ , which can be considered hydrogen-like with accuracy sufficient for our initial calculations. Use was made of the relevant theory [17] with exact computation of the integrals appearing in the relations derived in [17], in our calculations of  $\kappa$ , and that, by the way, leads to the difference of about one whole order of magnitude in the approximate expressions of those integrals cited in this paper. The value of  $\sigma$  was assigned as in Eq. (1.5). The equilibrium ionization (directly downstream of the evaporation wave) was determined from Saha's equation, using exact statistical weights and the energy levels of a nonhydrogen-like atom. In the calculations of the nonequilibrium absorption coefficient, the degree of ionization was determined from the kinetics equations cited earlier, and the population of the levels was assumed Boltzmannian, but corresponding to the electron temperature  $T_e$ .

The rate constant for ionization of the atom from the ground state by electron impact [16] is employed as  $\alpha_p$  in Eq. (1.2). Then  $\beta_p$  was determined such that the correct degree of ionization  $\alpha$  for the stated value of  $T_e$  would follow from Eq. (1.2) at ionization equilibrium.

We can assume the value of  $dv/vdt$  to be specified, and find it from the solution of the problem dealing with adiabatic disintegration of the gas in a centered rarefaction wave propagating in the wake of the evaporation wave. The assumption that the density can be found from the adiabatic solution is quite justified up to the onset of the flash, since the flash sets in, even according to the equilibrium theory, when the optical thickness of the layer of vapor is well below unity, and treatment of nonequilibrium shifts the onset of the flash even further into the range of lesser vapor thicknesses and shorter times, i.e., the amount of energy released is small. We note that a comparison was made, for the case of a purely equilibrium flash, between the criterion for the onset of the development of the flash obtained under the assumption that the term  $dv/vdt$  is specified while the rarefaction wave is adiabatic, and results of direct calculations in the corresponding gas-dynamics problem when absorption in the vapor, i.e., nonadiabaticity, is taken into account. The amount of agreement was found to be entirely satisfactory.

The rarefaction wave is centered since the flux density of the incident radiation is assumed constant below, and heat transfer into the bulk of the solid can be neglected, whereupon the rate constant of the propagation of the evaporation wave is also constant, as are all of the parameters in the transparent vapor directly downstream of the evaporation wave. The greatest inaccuracy here is due to the assumption that heat transfer to the bulk phase below the surface is absent, since we realize from the results of the calculations that the development times of shielding are commensurate with, or even less than, the transient periods of the stationary propagation rate of the evaporation wave (here we must recall that all of the times are reckoned in what follows from the instant evaporation commences).

It was assumed that  $T_e = T_i$  directly behind the evaporation wave, and the degree of ionization is equilibrium. In our work, calculations of parameters on the evaporation wave were carried out on the basis of the assumption that the evaporation temperature  $T_w$  downstream of the wave is equal to the temperature  $T_v$  of phase equilibrium between the liquid material and its single-atom vapor. We then determine  $T_v$  from the pressure downstream of the evaporation wave,  $p_w$ , which is lower than the pressure  $p_0$  on the surface where the evaporation commences, by roughly  $(\gamma + 1)$  times. That is the case if we assume that continuous

quasi-equilibrium "volume" evaporation takes place within the transition layer, while separate bubbles, and then a "droplet sheet," show up precisely at the beginning, while the droplets disappear gradually [4, 9, 10]. The same state of affairs is encountered in the case where the equilibrium surface evaporation [2] or non-equilibrium surface evaporation [1, 18-20] prevails, but equilibrium surface evaporation becomes predominant by the end of the transition region [9, 10, 14, 18].

If surface evaporation takes place and there is not enough time for equilibrium to become established between the vapor and the liquid, then the parameters downstream of the evaporation wave must be determined from an analysis of the structure of the zone on a level of molecular kinetics. An approximate treatment of that problem can be found in [18, 19].

This paper neglects distortion of the distribution function of the molecules moving toward the surface, because molecules moving away from the surface at first become scattered into that group; it is assumed that relaxation to a new equilibrium function occurs at distances of several free path lengths of the molecules, and the heat flux directed away from the body and due to ordinary heat conduction out of that zone, thereby broadening the transition zone, is not taken into further account at that point either.

According to published results [18, 19], the vapor temperature downstream of the evaporation wave is found to be lower than in the case where complete equilibrium has been established. That leads to lower degrees of ionization at the surface, and to later times at which screening takes place (later than in the assumption of phase equilibrium). However, since the vapor is supersaturated in such a situation, condensation takes place within comparatively short time periods according to estimates [14, 20] (quenching occurs only at points remote from the surface). Establishment of phase equilibrium therefore occurs at short distances from the surface, and we are free to make use of the relationship between temperature and pressure which we in fact adopted in our work:  $T_W = T_V(p_W)$ .

Evaporation is assumed to be surface evaporation in [2, 11], but it is also assumed that, because of the pronounced conduction of heat in the transition zone, the value of  $T_W$  is approximately equal to the temperature  $T_V$  of the phase transition at pressures  $p_0$  on the evaporating surface itself, i.e., it is  $(\gamma + 1)$  times higher than assumed by us. As those estimates indicate, the upshot would be a time  $t_*$  only one-third as long.

It is clear from the above that  $t_*$  is sensitive to the assumptions entertained as to the nature of evaporation and as to the structure of the transition zone. At the same time, the mechanical and thermodynamic parameters such as  $p$ ,  $u$ ,  $T$  are weakly dependent upon those assumptions, and the accuracy of current experiments is not adequate for a definitive answer to the question of what mechanism is involved in the evaporation [3, 11]. Comparison of theoretical research findings on the process of development of screening and experimental data might therefore provide some subsequent answer to this most interesting question.

If the relationship between  $T_W$  and  $p_W$  is known, and if it is also established that evaporation has gone to completion, then the relationship between the pressure  $p_W$  and the flux density  $q$  will be found from the conservation laws, and from the Jouguet conditions for constant flux and no conduction of heat. We also have to know the reflection coefficient  $k_r$ . According to theory, that coefficient would change with melting and with a further rise in temperature, whereupon it would contrast strongly with its "tabular" value in the cold state. Moreover, a change in the smoothness of the surface upon evaporation, the appearance of a droplet sheet in volume evaporation, can also lead to a change in that coefficient. The experimental data [20, 21] show that a twofold to threefold drop in the reflection coefficient (approximately) can occur. There are no direct experimental data on  $k_r$  for the substances discussed here. That introduces a corresponding uncertainty into the subsequent calculations. But it is evident that  $k_r$  seriously affects the relationship between the pressure  $p_0$  and the flux density  $q$ . The value of the effective reflection coefficient was therefore assigned such that measurements of the relationship between the momentum  $J$  and the energy  $E$  [3, 7] would be satisfied.

As an illustrative example, we cite the results of calculations for Al. The reflection coefficient was assigned the value 0.72, which leads to the ratio  $J/E = 4 \text{ dynes} \cdot \text{sec}/J$  for the unscreened mode of evaporation. The elastic scattering cross section  $\sigma$  was assigned the value  $20.4\pi a_0^2$ . The dependence  $T_V(p)$  was determined out to  $p = 110 \text{ atm}$  on the basis of [22], and at the higher pressures encountered in our examples by means of extrapolating those data.

TABLE 1.

$p_w$ , bar	$T_w$ , °K	$q_w$ , MW/cm <sup>2</sup>	$x_i^w$ , cm <sup>2</sup> /g	$x_n^w$ , cm <sup>2</sup> /g	$p_s$ , bar
75.4	4100	53.0	2.36	1.01	48.9
152.0	4500	105.0	7.41	3.37	69.4
274.0	4900	186.0	19.4	8.72	91.2
452.0	5300	304.0	44.2	20.3	114.0

TABLE 2.

$\epsilon = 1.78$ eV			$\epsilon = 1.16$ eV		
$q$	$t_*^v$	$t_*^s$	$q$	$t_*^v$	$t_*^s$
43.5	3.00	1.80	43.5	1.50	1.10
48.5	2.00	1.30	48.5	1.00	0.70
53.0	1.40	0.90	53.0	0.66	0.50
58.0	1.10	0.70	58.0	0.45	0.33

TABLE 3.

$\epsilon = 1.78$ eV			$\epsilon = 1.16$ eV		
$q$	$t_*(1)$	$t_*(2)$	$q$	$t_*(1)$	$t_*(2)$
53	$0.14 \cdot 10^1$	$0.14 \cdot 10^1$	53	0.66	0.61
105	0.16	0.14	105	$0.72 \cdot 10^{-1}$	$0.61 \cdot 10^{-1}$
186	$0.24 \cdot 10^{-1}$	$0.18 \cdot 10^{-1}$	186	$0.11 \cdot 10^{-1}$	$0.88 \cdot 10^{-2}$
303	$0.51 \cdot 10^{-2}$	$0.44 \cdot 10^{-2}$	303	$0.32 \cdot 10^{-2}$	$0.24 \cdot 10^{-2}$

Table 1 lists values of all of the parameters downstream of the evaporation wave (indicated by subscript w) and the pressure  $p_s$  on the shock wave which propagates through air, as needed for calculating the development of the flare in air (those calculations will be described in the next section).

The results obtained from solving Eqs. (1.1) and (1.2) for Al when the energies of the incident photons  $\epsilon = 1.78$  eV (1) (in the case of a ruby laser) and  $\epsilon = 1.16$  eV (2) (neodymium-doped glass laser) are plotted in Fig. 1 as a relationship of the time  $t$  ( $\mu$ sec) required for the flare to develop (after evaporation has set in) and the radiation flux density (in MW/cm<sup>2</sup>) under the assumption that complete equilibrium has been attained (broken curves), and with the difference between  $T_e$  and  $T_i$  taken into account (continuous curves). It is readily seen that treatment of that effect seriously lowers the "critical" density  $q_*$  corresponding to the onset of screening within a specified time  $t$  or the time  $t_*$  that it takes screening to develop at a specified flux density  $q$ .

Figure 2 shows the dependence of the behavior of solutions of the system on the time of arrival of a particle traversing the evaporation wave. The example taken is Al, flux density 186 MW/cm<sup>2</sup>, photon energy  $\epsilon = 1.78$  eV (behavior of solutions similar for other elements and other flux densities). The dashed curves represent  $T_e$ , the continuous curves represent  $T_i$ , the temperatures  $T_e$  and  $T_i$  are plotted in eV units, the time  $t$  in  $\mu$ sec. The time of the flare is plotted as abscissa  $t_p$  on the basis of equilibrium theory [9, 10], while  $t_{ne}$  is plotted on the basis of the nonequilibrium theory elaborated above. If integration is begun from a sufficiently short time (i.e., when the adiabatic cooling term is very small at the initial instant), then no flares can occur (see curves 1). As that time grows longer, the adiabatic cooling term becomes smaller, the relationship between it and the absorption term changes, so that  $T_e$  and  $T_i$  begin to rise rapidly at some time (curves 2). The time when  $\alpha = 0.4$  is taken arbitrarily as the time of the flare.

Figure 3 shows the same dependence for Al, but at a lower flux (53 MW/cm<sup>2</sup>). Comparison with Fig. 2 shows that the curves sharply alter the nature of their behavior as a function of the onset of integration. The difference  $\Delta = T_e - T_i$  contracts, and the flare time  $t_{ne}$  comes closer to  $t_p$ .

Figure 4 displays the distribution of the ion temperature and electron temperature with respect to the coordinate  $m/m_w$  at different times reckoned from the onset of evaporation.

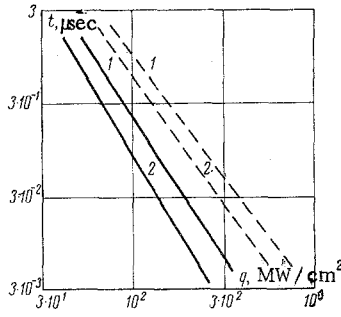


Fig. 1

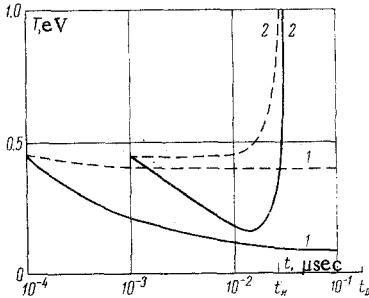


Fig. 2

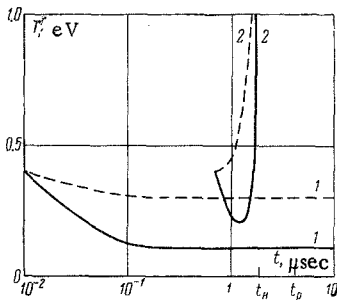


Fig. 3

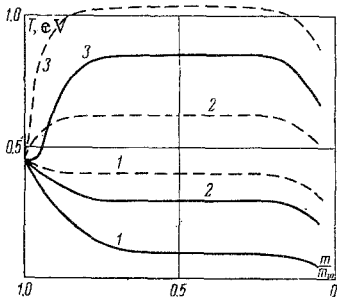


Fig. 4

The mass is referred to the total mass  $m_w$  of the evaporated material (Al). The flux density is  $186 \text{ MW/cm}^2$ , and  $\epsilon = 1.78 \text{ eV}$ . Curves marked by the numerals 1, 2, and 3 refer to the times  $5 \cdot 10^{-3}$ ,  $10^{-2}$ , and  $2.4 \cdot 10^{-2} \mu\text{sec}$ , respectively. The topmost curves depict the electron temperature distribution, while the bottom curves depict the ion temperature distribution. It is clear from the diagrams plotted that not only is there a change in the flare time when the deviation of the electron temperature from the ion temperature is taken into cognizance, but the entire temperature distribution pattern experiences a conspicuous change. In the equilibrium case,  $T_e$  and  $T_i$  both decline rapidly with increasing distance from the evaporating surface, and the flare takes place not far from the surface. When the departure from equilibrium is taken into consideration, we find that heating of the electrons is significant, and occurs at a great distance from the surface. That seems to be due to the increased relaxation time as the density of the material diminishes. The variable  $\kappa$  increases with increasing  $T_e$ , and the screening layer shapes up closer to the vacuum interface.

2. Similar estimating calculations of the time of onset of the sudden absorption increase were carried out, but for the case where the vapor expands and disintegrates into the surroundings. The ambient medium is assumed transparent to the radiation impinging on the surface of the evaporating body, because the mechanism underlying the increase in vapor energy does not differ from the case of expansion and disintegration into a vacuum, if the vapor pressure across the evaporation wave is much greater than the ambient pressure. However, when that condition is violated, the surroundings begin to hinder the expansion of the vapor considerably, so that energy losses in adiabatic expansion are cut down. As a consequence, the conditions for flare are satisfied at an earlier point than in the point of expansion and disintegration into a vacuum.

The development of the process of adiabatic disintegration and expansion of the vapor into the surroundings is depicted schematically in Fig. 5. The numeral 1 denotes the evaporation wave, 2 denotes one of the characteristics in the centered rarefaction wave, 3 denotes the contact discontinuity, 4 denotes the shock wave propagating through the medium. Figure 5 shows the pressure distribution at the instant of time  $t^*$ . In the layer where the pressure is equalized, heating proceeds from the onset of the radiation (the critical nature of the flare vanishes). The highest temperatures will be attained in those particles which evaporated ahead of all the others, i.e., the flare occurs in layers of the material which are adjacent to the surrounding medium.

Calculations of the kinetic equations of the process by which the flare of the vapor expanding into the surroundings develops were carried out by the same program, but leaving the surroundings out of account. The changes applied to the term accounting for the work done by forces of adiabatic expansion, and precisely at  $t_0 \leq t \leq t_k$  the term is considered the same as in the case of expansion into the vacuum, while it is assumed to be zero at  $t > t_k$ .

Here  $t_0$  is the time of onset of expansion of that element of the gas, and  $t_k$  is the time the expansion of that particle terminates, i.e., the time when the pressure in that particle, initially equal to  $p_w$ , falls to  $p_s$ . The results of the corresponding calculations are listed in Table 2. The superscript  $v$  indicates the value of  $t_*$  for the case of expansion into the vacuum, and the superscript  $s$  applies to the case of expansion into air of normal density ( $t_*^v$  and  $t_*^s$  are stated in  $\mu\text{sec}$ ). It is clear from the tabular data that the resistance presented by the medium shortens the time it takes the flare to develop as compared to the case of expansion into a vacuum, i.e., facilitates the formation of plasma.

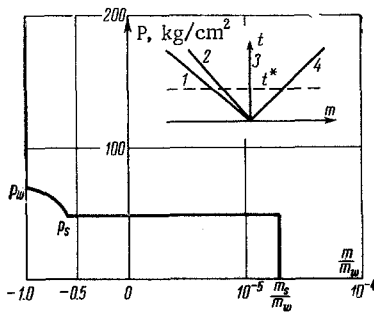


Fig. 5

3. The development of the flare takes place in the vapor, which is weakly ionized ( $\alpha \approx 10^{-3}$ ), but has a high density ( $\rho \approx 10^{-2}$  g/cm<sup>3</sup>). The estimates show that, throughout the process of development of the flare, the state of the vapor departs from the ideal state (i.e., the energy of the Coulomb interaction of charged particles becomes commensurate with the energy of random thermal motion). The Coulomb interactions affect the state of the gas in two ways, diminish the energy and pressure and, what is even more essential, shift the ionization equilibrium in the direction of larger  $\alpha$  values, increasing  $\alpha$  and thereby decreasing  $t_*$ .

The correction applied to the first ionization potential  $\Delta I_1$  under the assumption  $\Delta I_1/T_e \ll 1$  (this condition is fulfilled in the case of radiation of flux density  $q < 300$  MW/cm<sup>2</sup>) exhibits the form [16]

$$\Delta I_1 = 2c^3 \left( \frac{2\pi n_e}{k T_e} \right)^{1/2} \quad (3.1)$$

Since the correction  $\Delta I_1$  is a function of  $\alpha$ , the resulting degree of ionization determines a new correction to the ionization potential which leads anew to a change in the degree of ionization, and so forth. Saha's equations must be solved by the method of successive iterations in order to arrive at an exact determination of the degree of ionization and of the correction  $\Delta I_1$ . Within its range of validity (3.1), the iterative process converges rapidly. The results of the appropriate calculations are entered in Table 3. The superscript (1) denotes a value of  $t_*$  computed while ignoring departures from ideal behavior, while the superscript (2) denotes the same with the shift in ionization potential taken into account. Comparison shows that treatment of departures from the ideal state brings about a 15-20% reduction in the flare time for  $q \approx 150-300$  MW/cm<sup>2</sup>, while at lesser flux densities the effect of nonideal behavior on the flare time is less pronounced, and at flux densities  $q < 50$  MW/cm<sup>2</sup> the effect of departures from ideal behavior on flare time can be safely neglected. In the range of flux densities  $q > 300$  MW/cm<sup>2</sup>, we find an abrupt increase in  $\alpha$  due to nonideal behavior. That apparently constitutes evidence that there is a time when it is in general not possible to introduce the concept of a transparent evaporation wave or subsequent screening of that wave.

4. Comparison of calculations and experiment [12], where the duration of a ruby laser pulse (at the half-power level) was of the order of  $3.5 \cdot 10^{-2}$   $\mu$ sec (spoiled-Q mode), shows that the shielding times are in agreement in order of magnitude (calculated value  $q_* = 140$  MW/cm<sup>2</sup>, with experiment [12] showing that the plasma was heated up to  $T > 1$  eV at a mean flux density  $q > 75$  MW/cm<sup>2</sup>).

It should be borne in mind that the shape of the pulse was assumed to be rectangular in the calculations, hydrodynamic expansion was considered given, and the effect of heat conduction was left out of account, so that some discrepancy between the calculations and experiment [12] is hardly surprising. It is possible, however, that that could be due to some other physical effects that escaped attention, or to some inaccuracy in the values of the constants used. In any case, agreement between theory and experiment is far better than in the case where no account is taken of departures from equilibrium.

The maximum of the ratio of the mechanical momentum  $J$  to the energy supplied  $E$  was observed in [7] at the point where the duration of the ruby laser action was  $\approx 10^{-2}$   $\mu$ sec, and that corresponded to flux densities  $q$  of the order of  $\approx 500-1000$  MW/cm<sup>2</sup>. That has been related [11] to the onset of screening. According to calculations, the flare occurs long before that (at  $q = 200$  MW/cm<sup>2</sup>). A slight increase in the  $J/E$  ratio after the onset of the flare can be ascribed to cessation of heat transfer to the bulk of the material, a lessening of reflections, and of heat losses by evaporation (all of the radiation becomes absorbed by the plasma). A further decrease in the  $J/E$  ratio with increasing  $q$  is related to the rise in the temperature attained, and consequently to the specific energy per unit mass of vapor heated. That can be demonstrated more precisely only in the results of calculations of the "post-flare" stage of the procedure followed in [10], with heat transfer from the surface taken into cognizance (and additional evaporation after screening because of energy supplied from layers of the metal that have already become heated also taken into account), and by considering the real shape of the pulse, and the tabulated equation of state in the region of multiple ionization. The results of those calculations, and comparison of those results with measurements [12] taken of the pressure on the surface and the maximum brightness temperature as a function of the time, are to be published separately.



The behavior of the J/E ratio as a function of q for the time a neodymium-doped glass laser is active  $\tau = 3.5 \cdot 10^{-2} \mu\text{sec}$  has been reported in [20]. Here the peak J/E ratio lies within the range from 300 to 700 MW/cm<sup>2</sup>. A decrease in the critical flux density  $q_*$  with increasing  $\tau$  and with the transition from higher photon energies to lower photon energies is in accordance with the theory outlined. The theoretically predicted value of  $q \approx 120 \text{ MW/cm}^2$ .

An experimental investigation has been made [3] of free-running laser action under "weak" focusing conditions with the duration of a single flare  $\approx 1 \mu\text{sec}$  and characteristic "spot" dimensions  $r = 0.3 \text{ cm}$ . The conditions for the flatness of the disintegration, within the confines of a single spike, were observed. The flux density averaged over a pulse was  $4 \text{ MW/cm}^2$ , while the average taken over a single spike was  $18 \text{ MW/cm}^2$  (at duty ratio 0.22). Accordingly, the experimental conditions in [3] were close to those at which Al vapors are heated to higher temperatures than the phase-transition temperature, according to the computed data (the flux density was only one-third the calculated critical flux density). Results of calculations for other elements have been reported [13]. For example, the flux density required for screening to develop is far less in the case of W and Ta (not investigated in [3]) than in the case of Al, viz.  $\approx 10 \text{ MW/cm}^2$ , and that flux density can be observed in experiments with the laser parameters employed in [3]. It would be of interest to stage such verifying experiments.

Stronger focusing ( $r_0 \approx 0.03 \text{ cm}$ ) has been resorted to [1, 23]. The mean flux density attained a level of  $40 \text{ MW/cm}^2$ ; the average taken over a single spike was already  $180 \text{ MW/cm}^2$ . At that value of q, the theoretically predicted time  $t_*$  is  $10^{-2} \mu\text{sec}$ , and during that time the disintegration pattern is close to a flat one. It seems, on that account, that a heated plasma must be generated under those experimental conditions. But direct experimental measurements of the parameters characterizing the vapor at the bottom of the crater (e.g., the rates of disintegration) are lacking. One plausible indirect proof is the fact that the depth of the crater remains the same in the transition from the average (taken over the pulse) flux density  $q = 24 \text{ MW/cm}^2$  to the density  $q = 40 \text{ MW/cm}^2$ , or from 50 to  $160 \text{ MW/cm}^2$  on the average taken over a spike.

The comparison with experiment is a tentative one. The specific reason for that is, as alluded to earlier, the use of several simplifying assumptions in the exposition. For example, equilibrium states of the population of excited levels were assumed in calculating the absorption coefficient, because of bound-free transitions (one reference [24] indicates a possibility of the population deviating from equilibrium, and thereby leading to a slight change in absorption coefficient under the conditions in [24]); possible molecular absorption was ignored in calculating the absorption coefficients (according to the equilibrium theory, several percent Al<sub>2</sub> molecules are retained in the vapor at pressures of the order of 10 to 1000 bar, on the phase equilibrium curve), and also in calculating absorption by the wings of the broadened lines. In addition, departures of the evaporation process from equilibrium may have a substantial effect, as may the existence of surface ionization or the presence of any other physical processes at work. It is proposed that calculations in which the possible existence of some of those processes is taken into account to the maximum degree possible be carried out in subsequent work, so that a more detailed comparison with experiment can be performed. But practically all of the simplifying assumptions entertained lead to a still earlier onset of the screening phenomenon, which implies the possibility of laser generation of plasma in the range of very low flux densities, and consequently at fairly large dimensions of the irradiated spot and plasma cloud, and much greater lifetimes of the plasma cloud.

This will facilitate a whole series of physics and physical gas-dynamics research projects.

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