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OPTICAL BREAKDOWN OF MERCURY VAPOR

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Experiments on laser breakdown of gases with high ionization potential and with low excitation-band width are satisfactorily explained by the cascade theory proposed by Ya. B. Zel'dovich and Yu. P. Raizer [1]. The definition of the threshold electric field intensity in the light wave contains the quantity α - the probability that the electron will jump through the excitation-loss band - which is determined by solving the average stationary kinetic equation for a level with a certain average energy $\bar{\epsilon}$ in the excitation-loss band. However, in the case of atoms with small ionization potential and a broad excitation band, $\Delta \sim I_i$ (alkali metals, mercury [5]), an estimate of the probability α obviously calls for a detailed analysis of the energy levels of the atoms and for a solution of the complete quantum-kinetic equation in each band that is determined by the arrangement of the energy levels. For example, in the spectrum of the mercury atoms we can separate four bands: band I corresponds to the energy interval from the ground state level (I_0) to the first excited one (I_1^*), band II corresponds to the interval between levels 6^3P_1 and 6^1P_1, band III to the interval between levels 6^1P_1 and 7^1P_1, and band IV to the interval between the levels 7^1P_1 and the boundary of the continuum. The width of this band is $\Delta_4 = I_1 - I_3^* \sim \hbar\omega$, i.e., of the order of the ruby-laser quantum. This choice of the width of band IV is explained by the fact that the loss to the excitation level in this band will not influence the rate of development of the cascade, owing to the large probability of one-photon ionization.

The quantum-kinetic equation for the distribution function of the electrons $n(\epsilon)$, under the same assumptions as in [1], can be written in each band in the form

$$2\epsilon u \frac{d^2 n_k}{d\epsilon^2} + u \frac{d n_k}{d\epsilon} - \frac{n_k}{\theta_k} = 0, \qquad (1)$$

where k = 1, 2, 3, and 4 are the numbers of the bands, u is the "macroscopic" particle velocity along the energy axis, θ_1 is the time constant of the cascade development: $n(t) = n_0 \exp(t/\theta_1)$,

$$\frac{1}{\theta_2} = \frac{1}{\theta_1} + \frac{1}{\tau_1^*}$$

(τ_1^* - lifetime of the electron with respect to the excitation of the level I_1^*),

$$\frac{1}{\theta_3} = \frac{1}{\theta_1} + \frac{1}{\tau_{2i}^*} + \frac{1}{\tau_{21}^*}$$

(τ_{21}^* and τ_2^* - lifetimes of the electrons with respect to the excitation of the levels I_1^* and I_2^*, respectively),

$$\frac{1}{\theta_4} = \frac{1}{\theta_1} + \frac{1}{\tau_3^*} + \frac{1}{\tau_{31}^*} + \frac{1}{\tau_{32}^*}$$

(τ_3^*, τ_{31}^*, and τ_{32}^* - lifetimes of the electrons with respect to excitation of the levels I_3^*, I_1^*, and I_2^*, respectively).

The solutions of these equations in each band take the form

$$n_k(\epsilon) = C_1^k \exp\left(\sqrt{\frac{2\epsilon}{u\theta_k}}\right) + C_2^k \exp\left(-\sqrt{\frac{2\epsilon}{u\theta_k}}\right). \tag{2}$$

By making these solutions continuous on the band boundaries, under the conditions

$$n_k(\epsilon)|_{\epsilon = I_k^*} = n_{k+1}(\epsilon)|_{\epsilon = I_k^*}, \; J_k(\epsilon)|_{\epsilon = I_k^*} = J_{k+1}(\epsilon)|_{\epsilon = I_k^*} \tag{3}$$

(where k = 1, 2, 3, and $J_k(\epsilon)$ is the flux along the energy axis, defined in the same manner as in [1]), and using the boundary conditions

$$\begin{aligned} &n(\epsilon)|_{\epsilon = I_1} = 0 \qquad (I_1 = I_1 + 1 - 3\ \text{eV}), \\ &J_1(0) = 2J_4(I_1) + J_1(I_1^*) - J_4(I_1) = J_1(I_1^*)(1+\alpha), \end{aligned} \tag{4}$$

we can obtain an expression for the time constant θ_1 of the cascade development, and for the probability of "jumping" through the excitation band $\alpha = J_4(I_1)/J_1(I_1^*)$:

$$\frac{J_1(0)}{J_1(I_1^*)} = 1 + \alpha = \frac{\operatorname{sh}\kappa}{\kappa}, \quad \kappa = \left(\frac{2I_1^*}{u\theta_1}\right)^{1/2}. \tag{5}$$

Further, assuming that $\alpha \ll 1$ and expanding sinh κ in a series, we get

$$\theta_1 = \frac{I_1^*}{3u\alpha}, \tag{6}$$

where

$$\alpha = \frac{J_4(I_1)}{J_1(I_1^*)} = 2\sqrt{\frac{I_1}{I_1^*}} \exp\left\{-\sum_k \frac{\Delta_k}{(2uI_k^*\theta_k)^{1/2}}\right\} \quad (k = 2,3,4). \tag{7}$$

From (6) we can readily obtain an expression for the threshold intensity of the electric field:

$$E_k^2 = \frac{I_1^* m\omega^2}{e^2 \nu_{\text{eff}}(1 + 2\hbar\omega/I_1)u\theta_k}. \tag{8}$$

If we take into account the fact that electrons having an energy I_1 larger by 1 - 3 eV than the ionization potential I_i not only ionize but also excite the atoms, then the boundary condition for the fluxes (4), with correction for the single-photon ionization from the upper levels, takes the form

$$J_1(0) = \frac{2J_4(I_1)}{1 + r_i/r_\Sigma^*} + J_4(I_1)\left(1 - \frac{1}{1 + r_i/r_\Sigma^*}\right) + \\ + J_1(I_1^*) - J_4(I_1) + \beta \frac{J_1(I_1^*) - J_4(I_1)}{1 + r_3^*/r_\Sigma^*}, \tag{9}$$

where τ_Σ^* is the lifetime of the atom with respect to the excitation, τ_i - lifetime of the atom with respect to ionization, β - ratio of the number of excited atoms with energy $I_k^* \geq I - \hbar\omega$ to the total number of excited atoms:

$$\beta = \frac{J_3(I_3^*) - J_4(I_1)}{J_1(I_1^*) - J_4(I_1)} = \frac{1/a_{34} - 1}{1/a - 1},$$

where

$$a_{34} = 2\sqrt{\frac{I_1}{I_3^*}} \cdot \exp\left\{ - \frac{\Delta_4}{(2I_1 \dot{v} \theta_4^*)^{1/2}} \right\} \tag{10}$$

is the probability that the electron will "jump" through the fourth band of width Δ_4.

Thus, allowance for single-photon ionization of the upper excited levels and for the energy lost by the electrons to excitation of the atoms in the energy interval $\epsilon > I_i$ leads to a replacement of α in (8) by α':

$$a' = \frac{a}{1 + r_i/r_\Sigma^*} + \frac{\beta(1 - a)}{1 + r_3^*/r_\Sigma^*}. \tag{11}$$

Figure 1 shows plots of the theoretical threshold curves calculated by means of the formulas from [1] and (8) with the correction (11) for mercury ($\Delta \sim I_i$) and xenon ($\Delta \ll I_i$). Agreement of the results of the calculation for xenon is to be expected, since formulas (7) and (8) go over into the corresponding formulas obtained by Ya. B. Zel'dovich and Yu. P. Raizer [1] if $\Delta \ll I_i$ and τ^* is constant for all the excitation bands. The difference between the results of both theories for mercury ($\Delta \sim I_i$) can be attributed to the fact that when account was taken of the photoionization of the upper levels of the excited atoms in the theory of [1] with the correction β from [2], it was assumed that the parameter $\beta = (\hbar\omega/\Delta)^2$ does not depend on the electric field intensity, whereas it is seen from (10) that β depens strongly on the field intensity, being large in strong fields (low pressures) and small in weak fields (high pressures). In addition, in the region of high pressures, exact allowance for the loss to excitation yields for α' lower values than in [1], leading

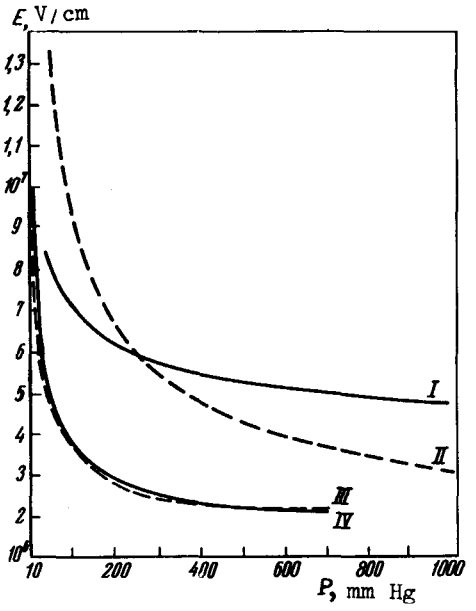

Fig. 1

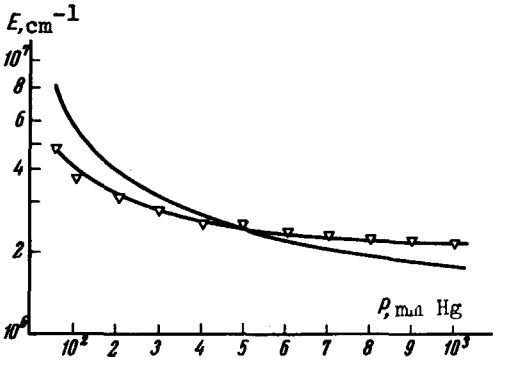

Fig. 2

to an increase of the threshold field.

To verify this theory, an experiment was performed in mercury vapor. We used a laser with Q-switching by a rotating total-internal-reflection prism. The pulse energy was ~1 J, the duration ~60 nsec, and the divergence ~6'. The laser radiation was focused by a lens with a focal distance f = 15 mm inside a glass vacuum chamber. To produce a definite mercury vapor pressure, the chamber was placed in an oven so constructed that the temperature gradient in the volume of the oven was minimal. The breakdown was recorded by a photoelectric method and was revealed by the appearance of mercury spectral lines in the spark emission spectrum.

Figure 2 shows the experimental and theoretical threshold plots against the mercury vapor pressure. The theoretical threshold field was estimated by means of formula (8) with allowance for the diffusion loss and for collisions of the second kind [3,4] and for a correction in accordance with formula (11). Figure 2 shows that the experimental and theoretical curves agree quite well.

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