charge and the inhomogeneity of $T_a(r)$. Second, a change in $T_e(r)$ leads to a much stronger change in $N_e(r)$ than in the equilibrium case. This apparently causes the very strong contraction of the current in pure argon. The transition of the current to the contracted state is connected with the fact that the diffusion mechanism of eliminating the charged particles gives way to the volume mechanism. The recombination rate is usually proportional to $N_e N_i$, and therefore the transition to the contracted state occurs first on the discharge axis, and a diffuse 'halo' remains on the periphery (see the figure), with a characteristic dimension that decreases with increasing current.

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LOSS MECHANISM IN PRIMARY OPTICAL BREAKDOWN

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Primary optical breakdown of a gas in the focus of a laser at pressures on the order of atmospheric can be satisfactorily explained on the basis of the electron-avalanche mechanism proposed by Ya. B. Zel'dovich and Yu. P. Raizer [1]. Within the framework of the avalanche theory it is also possible to account for processes that influence the development of the electron avalanche, such as photoionization of excited atoms [2] and diffusion losses [3]. In a number of cases the development of the avalanche may be greatly influenced also by impacts of the second kind between electrons and excited atoms, since these accelerate the avalanche. The time constant θ of avalanche development is given by

$$\theta^{-1} = \frac{\alpha}{\theta_0} - \frac{1}{\tau_D}, \tag{1}$$

where $\theta_0^{-1} = \dot{\varepsilon}/I_1$

$$\dot{\varepsilon} = \dot{\varepsilon}_0 - \dot{\varepsilon}_1 = \frac{e^2 E^2 \nu_e}{m(\omega^2 + \nu_e^2)} - \frac{2m}{M}\varepsilon\nu_e$$

is the rate of growth of the energy ε of the electron as a result of deceleration absorption of light quanta $h\omega$ in collisions between the electrons and the neutral atoms, I_1 is the effective ionization threshold, and α is the probability that the electron jumping through an excitation-loss band of width $\Delta = I_1 - \varepsilon^*$ (ε^* is the energy of the first excited level). The procedure for accounting for the photoionization of the upper levels, proposed in [2], leads to the substitution $\alpha' = \alpha + \beta(1 - \alpha)$, where β is the probability calculated by means of formula (7.2) of [1]. To take impacts of the second kind into account, we introduce the parameter $\gamma = W_{II}/W_I$ into account, characterizing the ratio of the probabilities of impacts of the second and first kinds. As a result of a collision of the second kind with an excited atom, the electron acquires an energy that is either sufficient for subsequent ionization of the

atom, or else an energy lower than I_1, but lying within the excitation band $\Delta = I_1 - \varepsilon$.* It is usually possible to assume that $\Delta \ll I_1$ and therefore, as shown by estimates, an electron experiencing an impact of the second kind acquires an energy equal to I_1 within a time much shorter than θ. Thus it can be assumed approximately that practically all impacts of the second kind lead to ionization of the atom. In this case the modified second boundary condition of Eq. (6.1) of [1] takes the form

$$J(0) = J(I_1)[2\alpha' + 2\gamma(1-\alpha') + (1-\alpha')(1-\gamma)] \tag{2}$$

and the allowance for the impacts of the second kind reduces to replacing α in (1) by α'':

$$\alpha'' = \alpha' + \gamma(1-\alpha'). \tag{3}$$

The parameter γ can be determined from the balance of exciting and de-exciting collisions, using the principle of detailed balancing, which relates the cross sections of impacts of the first and second kind under the assumption that the number of electrons increases exponentially in the focusing volume. Averaging the cross sections over the velocity interval, we get:

$$\gamma = 1 - \exp\left[\theta(1-e^{t/\theta})\frac{\langle v\sigma_{II}(v)\rangle}{V}\right] \tag{4}$$

(V is the focusing volume). It is easy to see that $\gamma = \gamma(t)$ has a strong dependence on the time, i.e., we can expect $\gamma \approx 0$ ($\alpha'' = \alpha'$) when $t \leq t_1$, and $\gamma = 1$ ($\alpha'' = 1$) when $t \geq t_1 + \Delta t$, with $\Delta t \ll \tau_u$ (τ_u is the laser pulse duration). After averaging we get

$$\overline{\alpha''} = 1 - \delta(1-\alpha'),\ \delta = t_1/\tau_u. \tag{5}$$

Replacing α by $\overline{\alpha''}$ in (1), we rewrite the latter in the form

$$1/\theta = 1/\theta_{00} - 1/\theta_{01} - 1/\tau_D - (1-\alpha')/\theta_0 + \frac{(1-\delta)(1-\alpha')}{\theta_0}. \tag{6}$$

Here $1/\theta_{00} = \dot{\varepsilon}_0/I_1$ is the time constant for electron heating in the field of the light wave, $1/\theta_{01} = \dot{\varepsilon}_1/I_1$ is a term characterizing the elastic losses, $1/\tau_D$ represents the diffusion loss, $(1-\alpha')/\theta_0$ represents the inelastic excitation loss, and $[(1-\alpha')(1-\delta)]/\theta_0$ is the term that takes into account the contribution made to the avalanche by the impacts of the second kind. Using the breakdown condition

$$1/\theta \equiv 1/\theta_k = \frac{\ln n(\tau_u)}{\tau_u}, \tag{7}$$

we can determine from (6) the threshold value of the electric field intensity in the light wave.

Figure 1 illustrates the relative role of the different processes that govern the time constant of the avalanche. The curves were plotted for argon under the following conditions: laser pulse duration $\tau_u = 80$ nsec, focusing-region dimension $r = 5.4 \times 10^{-3}$ cm. It is seen from the diagram that at low pressure an important role is played by the diffusion losses. At high pressures, to the contrary, the elastic and inelastic losses come into play, but the

inelastic losses are offset to a considerable degree by the impacts of the second kind. Figures 2 and 3 show the results of experimental measurements of the threshold electric-field intensities in Ar at pressures 360 - 700 mm Hg (Fig. 2) as functions of the dimension of the focusing region (laser pulse duration ∿ 65 nsec, laser beam divergence $\phi \sim 5'$) and at pressures 0.2 - 150 atm (Fig. 3, laser pulse duration 80 nsec, dimension of focusing region ∿ 5.4×10^{-3} cm). The measurements were made with a setup similar to that described in [3].

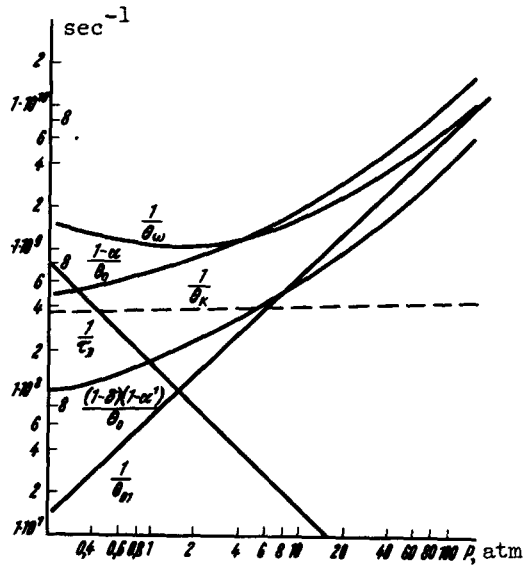

Fig. 1

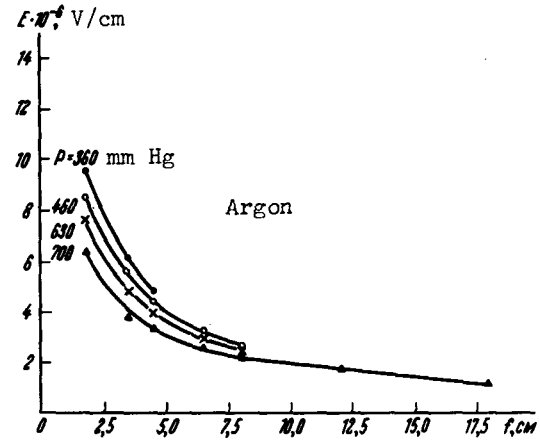

Fig. 2

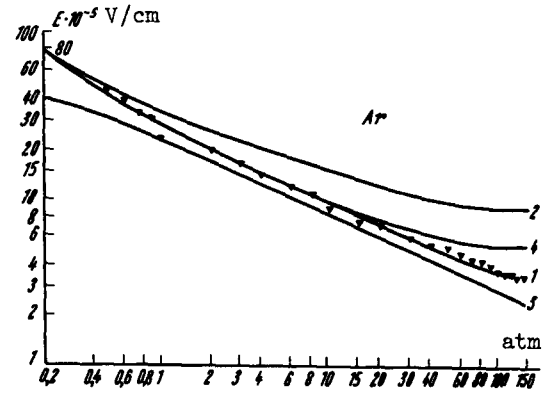

Fig. 3. 1 - experimental curve, 2, 3, 4 - theoretical, calculated with allowance for elastic losses (2), impacts of 2-nd kind (3), and both(4)

The table lists the ratios E_{theor}/E_{exp}, the theoretical values being calculated in accordance with the formulas of [1] (first column), with correction for the photoionization of the excited atoms [2] (second column, and in accord with formula (6) with allowance for the elastic and diffusion losses and impacts of the second kind (third column). The experimental

	I		II		III		Gas	Ar		Kr	
Gas / P, mmHg	Ar	He	Ar	He	Ar	He	f mm / P mmHg	18	65	18	65
1500	0.71	1.12	0.545	0.98	0.77	0.95	100	1.51	1.04	1.43	1.35
10^4	1.31	1.53	0.69	1.07	1.05	1.02	360	0,82	0.97	0.9	1.04
10^5	1.96	2.15	0.52	0.89	1.15	1.1	560	0.7	0,96	0.8	0.99

values E_{exp} were taken from the paper of Meyerand and Haught [4], from [3], and from the present measurements. It should be noted that the diffusion losses under the conditions of the experiments in [3] can be estimated by assuming a free diffusion mode. Indeed, the solution of the equation describing the time variation of the number of electrons (provided the diffusion coefficients do not depend on the energy), with allowance for the transient character of the diffusion [5], leads to a correction factor on the order of unity in the calculation of E_{theor} - the threshold electric field intensity.

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STIMULATED RESONANCE EFFECTS IN POTASSIUM VAPOR

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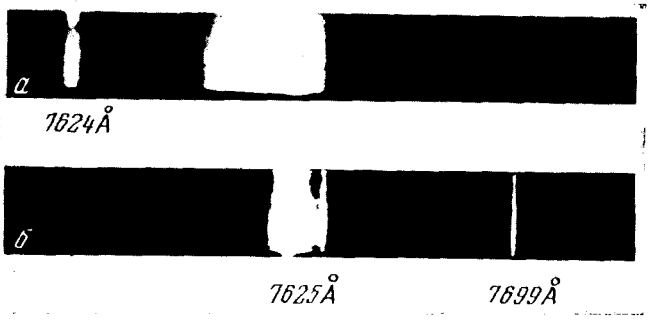
We present here experimental results on the passage of intense optical radiation through potassium vapor. We observed intense stimualted emission at wavelengths 7665 Å and 7699 Å coinciding with the resonant $4P_{3/2,1/2} - 4S_{1/2}$ transitions of the potassium atoms. Insofar as we know, no such effect has been reported in the literature. We also obtained stimulated emission at a wavelength 7624 Å as a result of Raman scattering from the excited electronic level $4P_{3/2}$ [1].

The intense-radiation source was a ruby laser with passive shutter, with an approximate power of 50 MW. A cell 10 cm long, filled with nitrobenzene and placed in the unfocused beam of the main radiation, produced an intense Stokes component of stimulated Raman scattering at a wavelength λ_2 = 7658 Å. The radiation with λ_1 = 6943 Å and λ_2 = 7658 Å was focused with a lens in a cuvette 10 cm long filled with potassium vapor. The cell with the potassium was heated to 250 - 350°C, corresponding to a vapor pressure 0.05 - 1.7 mm Hg. The light passing through the cell was registered with a DFS-13 spectrograph.

In one series of experiments, only light at wavelength λ_2 was focused on the cell, and the main ruby radiation at λ_1 was blocked with an FS-7 filter. In this case there appeared in the spectrum an intense component on the short-wave side, at 7624 Å (Fig. a). This radiation is connected with the two-photon stimulated effect, as a result of which the excited potassium atoms are transferred from the $4P_{3/2}$ level to $4P_{1/2}$. The same effect was observed in [1].



On the other hand, when the main laser