Light-induced current in sodium vapor

S. N. Atutov, I. M. Ermolaev, and A. M. Shalagin

Institute of Automation and Electronics, Siberian Division, USSR Academy of Sciences
(Submitted 10 November 1985)
Zh. Eksp. Teor. Fiz. **90**, 1963–1971 (June 1986)

A novel photovoltaic effect, light-induced current in a rarefied gas, is investigated experimentally and theoretically. An expression for the current is obtained for the case when the ions are produced by associative ionization of the atoms from the excited states. The influence of optical pumping is discussed. The experiments were performed in sodium vapor. A unique antisymmetric dependence of the current on the detuning from the radiation frequency is recorded for the first time. The dependences of the current on the resonance fluorescence, on the buffer-gas pressure, and on other factors are investigated. The experimental results agree well with the theory.

1. INTRODUCTION

A novel phenomenon, light-induced current in rarefied gases, was predicted and observed experimentally in Ref. 1. The gist of the phenomenon is that resonant action of the radiation on the rarefied gas can produce an electric current in the latter.

We recall briefly the physical nature of the effect. It is based on the fact that velocity-selective atom excitation causes all the excited particles to move in the same direction with velocities close to Ω/k, where Ω is the detuning of the radiation frequency from the center of the gas absorption line, and $\mathbf{k}$ is the wave vector. If such excited particles are ionized in some manner, both the electrons and ions produced thereby will have the same average velocity Ω/k. Since, however, the electron mass is much smaller than that of the ion, the absolute value of the electron velocity exceeds greatly the thermal velocity of the atoms and the directed-motion velocity Ω/k. In other words, the electrons travel practically isotropically and (in the absence of collisions) go off rapidly to the walls. Only ions remain in the volume and it is their directed motion which comprises the current. Depending on the sign of Ω, the current is either parallel or antiparallel to the wave vector. There is no current at exact resonance (at $\Omega = 0$).

The effect described is truly voltaic, in contrast to the known optovoltaic effect[2] that alters the conductivity of the gas (or plasma) and can occur only in the presence of a potential difference. On the other hand, similar optovoltaic effects were discussed in Refs. 3–5, and the physical basis of the effects considered both in these references and in Ref. 1 is close to the physical basis of the light-induced drift (LID) effect.[6] The effect considered in Ref. 1 and by us differs from those described in Refs. 3 and 4 in that decelerating bulk collisions are disregarded. It turns out that in the Knudsen situation, when the ion and electron momenta relax on the walls, the light-induced current can reach substantially higher values, and this has made it possible to detect it in experiment.

The effect was theoretically considered in Ref. 1 for the case of low densities of charged particles, when their interaction with the space charge can be neglected. The current calculated flows between plane-parallel electrodes in the space between which $Q(\mathbf{v})$ ions are produced per unit volume and per unit time in a unit velocity interval. The current was estimated in Ref. 1 under plasma conditions, i.e., when the Debye radius is much shorter than the distance d between the electrodes. This problem was examined in greater detail in Ref. 7, where the validity of using a hydrodynamic approach to the calculation of the current was justified. The corresponding hydrodynamic equations are of the form:

$$\begin{aligned} \frac{d}{dx}(nu) &= \int Q(\mathbf{v})\,d\mathbf{v}, \\ \frac{d}{dx}(nu^2+nv_s^2) &= \int Q(\mathbf{v})\,v_k\,d\mathbf{v}-\nu nu, \end{aligned} \tag{1}$$

where n and u are the density and directed velocity of the ions, v_S is the speed of sound in the plasma, ν is the frequency of the ion collisions that lead to the Maxwellian distribution, and the subscript k labels a component along the wave vector. Solving (1), we readily obtain an expression for the current:

$$J = \frac{ed}{2v_s}\frac{1}{1+\nu d/4v_s}\int Q(\mathbf{v})\,v_k\,d\mathbf{v}. \tag{2}$$

To calculate the current we must relate $Q(\mathbf{v})$ to the characteristics of the gas and of the radiation under the specific ionization conditions. This was done in Ref. 7 for the simplest case of a two-level system and for ionization from an excited state by supplementary radiation.

We investigated the light-induced current experimentally in sodium vapor, where no supplementary radiation is needed for the ionization, since associative ionization takes place from the excited state:

$$\mathrm{Na}(3P)+\mathrm{Na}(3P)\rightarrow \mathrm{Na}_2^{+}+e^{-}. \tag{3}$$

The onset of current in sodium vapor via the mechanism (3) is treated theoretically in Sec. 2. Account is taken there of the influence of the optical pumping due to the hyperfine splitting of the ground state.

In Sec. 3 we report the experimental results for the light-induced current in sodium vapor. We record for the

 0038-5646/86/061149-06$04.00

first time ever an antisymmetric dependence of the current on the detuning Ω. We investigate the dependences of the current on the resonant-fluorescence intensity and on the buffer gas pressure. The theory is compared with experiment.

2. THEORY

To calculate the light-induced current we must know the distribution functions of the excited atom in velocity and in the coordinates. In the simplest variant, the sodium atom can be modeled by a three-level system, with two levels (n and l) the components of the hyperfine splitting of the ground state $3S_{1/2}$, and the third (m) the excited state ($3P_{1/2}$ or $3P_{3/2}$).

We assume the pressure low and neglect the effect of collisions on the density matrix. In addition, we assume that the homogeneous linewidths are much smaller than the Doppler parameter $k\bar{v}$ and the spacing Δ of the hyperfine-structure components. In our case the field interacts independently with each of the components (the Bennett peaks do not overlap). The contributions from both transitions to the distribution of the excited atoms is additive under these conditions. Without loss of generality, we can therefore write the equations for the density matrix under the assumption that the field acts on only one of the transitions (for the sake of argument, m–n). We write down these equations in the reference frame of the atom passing through the light beam:

$$\begin{aligned}
&(d/dt+2\Gamma)\rho_m=-2\,\mathrm{Re}[iG_{mn}^*(t)\rho_{mn}],\\
&d\rho_n/dt-\Gamma\rho_m=2\,\mathrm{Re}[iG_{mn}^*(t)\rho_{mn}],\\
&[d/dt+\Gamma-i(\Omega-\mathbf{kv})]\rho_{mn}=iG_{mn}(t)(\rho_n-\rho_m),\\
&d\rho_l/dt-\Gamma\rho_m=0,\quad G_{mn}(t)=E(t)d_{mn}/2\hbar,\quad \Omega=\omega-\omega_{mn}.
\end{aligned}\tag{4}$$

Here $E(t)$ and ω are the amplitude of the electric field and the radiation frequency; d_{mn} and ω_{mn} are the dipole-moment matrix element and the frequency of the m–n transition; 2Γ is the decay constant of the excited state.

Equations (4) are written under the assumption that the frequency of the resonant transfer of the excitation is negligibly small compared with the spontaneous-decay rate Γ. We solve the system (4) with the initial conditions

$$\rho_{n,l}(t\to-\infty)=NB_{n,l}W(\mathbf{v}),\quad \rho_m(t\to-\infty)=0.$$

Here $B_{n,l}$ are Boltzmann factors. $W(\mathbf{v})$ is the equilibrium Maxwellian distribution, and N is the atom density.

We examine now the evolution of the density matrix as the atom travels through the light beam. The particle first enters the periphery of the beam, where the field is weak ($|G_{mn}|^2\ll\Gamma^2+(\Omega-\mathbf{kv})^2$). In this case we have as a solution of the system (4):

$$\begin{aligned}
&\rho_m(t)=B_nW(\mathbf{v})N\Lambda\varkappa_{mn}(t)\exp\Big[-\Gamma\Lambda\int_{-\infty}^{t}\varkappa_{mn}(t')\,dt'\Big],\\
&\rho_n(t)=B_nW(\mathbf{v})N\exp\Big[-\Gamma\Lambda\int_{-\infty}^{t}\varkappa_{mn}(t')\,dt'\Big],\\
&\rho_l(t)=W(\mathbf{v})N\Big\{1-B_n\exp\Big[-\Gamma\Lambda\int_{-\infty}^{t}\varkappa_{mn}(t')\,dt'\Big]\Big\},\\
&\Lambda=\Gamma^2/[\Gamma^2+(\Omega-\mathbf{kv})^2],\quad \varkappa_{mn}(t)=|G_{mn}(t)|^2/\Gamma^2.
\end{aligned}\tag{5}$$

Relations (5) describe optical transfer of resonant-velocity atoms from level n to level l via the excited state m. The transfer rate is characterized here by the argument of the exponential

$$\begin{aligned}
\Gamma\Lambda\int_{-\infty}^{t}\varkappa_{mn}(t')\,dt'&\sim\frac{\Lambda\Gamma\varkappa_{mn}^2}{2}\left(\frac{d\varkappa_{mn}}{dt}\right)^{-1}\\
&=\frac{\varkappa_{mn}\Lambda}{2}\left[\frac{d(\ln\varkappa_{mn})}{d(\Gamma t)}\right]^{-1}.
\end{aligned}\tag{6}$$

For allowed atomic transitions, the characteristic atomic time $1/\Gamma$ is much shorter than the time of passage of the particle through the light beam. Under these conditions the field increases slowly in the atom system (i.e., $d(\ln\varkappa_{mn})/d(\Gamma t)\ll1$). Thus, optical pumping can be effective in weak fields ($\varkappa_{mn}\Lambda\ll1$) because it proceeds quite slowly. After passing through the peripheral light beam and landing in the strong-field region (where $\varkappa_{mn}\Lambda\gtrsim1$), the atom turns out to be already optically pumped, i.e., in this region we have

$$\rho_n(t)=\rho_m(t)=0,\quad \rho_l(t)=W(\mathbf{v})N.$$

The behavior of the density matrix in optical fields of any intensity is thus well described by expressions (5), provided only that $\Gamma\gg\bar{v}/a$, where a is the beam radius.

We calculate first the current in the simplest case of optical pumping. These conditions are realized at a small saturation parameter ($\varkappa_{mn}\ll\bar{v}/\Gamma a$) and in the absorption-line wing ($\Omega\gg k\bar{v}$). We write in this case the expression for $\rho_m(r,\mathbf{v})$ with allowance for excitation from both hyperfine-structure components. Assuming the field to have the form of a Gaussian beam, $\varkappa_{mn}=\varkappa_{mn}^0\exp(-r^2/2a^2)$, we get

$$\begin{aligned}
\rho_m(r,\mathbf{v})=&B_nW(\mathbf{v})\frac{N\Gamma^2\varkappa_{mn}{}^0\exp(-r^2/2a^2)}{\Gamma^2+(\Omega'-\mathbf{kv}+\Delta/2)^2}\\
&+B_lW(\mathbf{v})\frac{N\Gamma^2\varkappa_{ml}{}^0\exp(-r^2/2a^2)}{\Gamma^2+(\Omega'-\mathbf{kv}-\Delta/2)^2}.
\end{aligned}\tag{7}$$

The number $Q(\mathbf{v})$ of ions produced per unit volume and per unit time in a unit velocity interval is calculated from the equation

$$Q(\mathbf{v})=\int\sigma(u)\,|\mathbf{u}|\,\rho_m(\mathbf{v}+\mathbf{u}/2)\rho_m(\mathbf{v}-\mathbf{u}/2)\,d\mathbf{u},\tag{8}$$

where $\sigma(u)$ is the associative-ionization cross section and $\mathbf{u}$ is the relative velocity of the colliding excited atoms.

In the region $\Omega\lesssim k\bar{v}$ the excitation is extremely velocity-selective, and it can be assumed in the calculation of the integral (8) that $\rho_m(\mathbf{v})$ is a δ-function of the velocity component v_k along the wave vector. Substituting (8) and (7) in (2) we obtain for an optically thin medium the following expression for the current (we assume for the time being that $vd/4v_S\ll1$)

$$J=\frac{\pi a^2 de}{2}\frac{\Gamma^2N^2}{(k\bar v)^2}\frac{\bar v}{v_s}\Big\{K(0)B_n^2(\varkappa_{mn}^0)^2\times\exp\Big[-2\Big(\frac{\Omega'+\Delta/2}{k\bar v}\Big)^2\Big]\times\frac{\Omega'+\Delta/2}{k\bar v}+2\frac{\Omega'}{k\bar v}K\Big(\frac{\Delta}{k}\Big)B_nB_l\varkappa_{mn}^0\varkappa_{ml}^0\exp[-2(\Omega'/k\bar v)^2]+K(0)B_l^2(\varkappa_{ml}^0)^2\exp\Big[-2\Big(\frac{\Omega'-\Delta/2}{k\bar v}\Big)^2\Big]\frac{\Omega'-\Delta/2}{k\bar v}\Big\},\quad(9)$$

$$K(u_k)=\frac{1}{\bar v^2}\int_0^\infty\exp\Big(-\frac{u^2}{2\bar v^2}\Big)\sigma((u_\perp^2+u_k^2)^{1/2})(u_\perp^2+u_k^2)^{1/2}u_\perp du_\perp.$$

Here $u_\perp$ is the orthogonal component of the velocity, and $K(u_k)$ has the meaning of the relaxation constant (3) under conditions when one of the components of the relative velocity of the particles is firmly specified and is equal to u_k, while the particle distribution over the remaining two velocity components is in equilibrium.

In the absorption line wing ($\Omega\gg k\bar v,\Delta$). When the excitation is practically balanced with respect to v_k, the current satisfies the relation

$$J=KN^2\frac{\pi a^2 de}{2}\frac{\bar v}{v_s}\frac{\Gamma^4k\bar v}{\Omega^5}(B_n\varkappa_{mn}^0+B_l\varkappa_{ml}^0)^2,\qquad K=\Big(\frac{2}{\pi}\Big)^{1/2}\frac{1}{\bar v^3}\int_0^\infty\exp\Big(-\frac{u^2}{2\bar v^2}\Big)\sigma(u)u^3du.\quad(10)$$

Here K is the reaction rate in a gas with equilibrium velocity distribution.

If we assume that $K(u_k)=K$, i.e., is independent of u_k, it is easy to obtain for the value of the light induced current an expression that is valid for any detuning Ω':

$$J=\frac{\pi a^2dN^2eK\bar v}{8\Gamma^2v_s}\Big\{p_n^2\varphi\Big(\Omega'+\frac{\Delta}{2}\Big)+p_np_l\Big[\varphi\Big(\Omega'+\frac{\Delta}{2}\Big)+\varphi\Big(\Omega'-\frac{\Delta}{2}\Big)\Big]+p_l^2\varphi\Big(\Omega'-\frac{\Delta}{2}\Big)\Big\}$$

$$p_{n,l}=2B_{n,l}\Gamma\varkappa_{mn,ml}^0\int_{-\infty}^{\infty}\frac{\Gamma^2W(v_k)dv_k}{\Gamma^2+(\Omega'-kv_k\pm\Delta/2)^2},$$

$$\varphi(\Omega)=\frac{\mathrm{Re}[zw(z)]}{\mathrm{Re}[w(z)]},\quad z=\frac{\Omega}{k\bar v}+i\frac{\Gamma}{k\bar v},\quad(11)$$

$$w(z)=e^{-z^2}\Big(1+\frac{2i}{\pi^{1/2}}\int_0^z e^{t^2}dt\Big).$$

Here $p_{n,l}$ is the probability of radiation absorption by an n- or l-level atom located on the beam axis, while $\varphi(\Omega)$ is a function introduced in Ref. 8 and describing the frequency dependence of the SID effect. Recognizing that the relation

$$p_{n,l}\varphi(\Omega'\pm\Delta/2)=-{}^1/_2k\bar v\,dp_{n,l}/d\Omega',$$

holds at low values of the saturation parameter, we can express the current in terms of radiation power $\Delta S(\Omega')$ absorbed in the cell:

$$J=-\frac{e\bar vK}{32\pi a^2dv_s}\frac{k\bar v}{2}\frac{d}{d\Omega'}\Big[\frac{\Delta S(\Omega')}{\hbar\omega\Gamma}\Big]^2.\quad(12)$$

Thus, when the optical pumping is negligible, the light-induced current line shape is the derivative of the square of the absorption-line profile. Since the hyperfine structure cannot be resolved in the $J(\Omega')$ plot.

We discuss now the case when the optical pump plays an essential role. To understand the ensuing qualitative changes, we turn to relations (5) and consider a simple exactly solvable case, when the field has the form of a one-dimensional Gaussian beam: $\varkappa_{mn}=\varkappa_{mn}^0\exp(-x^2/2a^2)$. We rewrite the expression for ρ_m from (5) in the laboratory frame:

$$\rho_m(x,\mathbf{v})=\begin{cases}NB_nW(\mathbf{v})\Lambda\varkappa_{mn}^0\exp\Big\{-\dfrac{\Gamma a}{v_x}\Lambda\displaystyle\int_{-\infty}^{x}\exp\Big(-\frac{x'^2}{2a^2}\Big)dx'\Big\}, & v_x>0,\\ NB_nW(\mathbf{v})\Lambda\varkappa_{mn}^0\exp\Big\{\dfrac{\Gamma a}{v_x}\Lambda\displaystyle\int_{x}^{\infty}\exp\Big(-\frac{x'^2}{2a^2}\Big)dx'\Big\}, & v_x<0.\end{cases}\quad(13)$$

Integrating (13) with respect to the coordinate, we obtain the excited-particle distribution function in velocity:

$$\rho_m(\mathbf{v})=B_nW(\mathbf{v})N\frac{|v_x|}{\Gamma a}a\Big[1-\exp\Big(-(2\pi)^{1/2}\frac{\Gamma a}{|v_x|}\varkappa_{mn}^0\Lambda\Big)\Big].\quad(14)$$

Figure 1 shows a plot of $\rho_m(v_k)$ at various saturation parameters. At small $\varkappa_{mn}^0$, when the argument of the exponential has an absolute value much smaller than unity, the optical pumping is insignificant and the plot of $\rho_m(v_k)$ has the usual Lorentz shape.

Let us note the properties of the Bennett peak under conditions when the optical pumping is strong. The amplitude of the peak is $\Gamma a/2|v_x|$ times smaller than the equilibrium distribution function of the particles on the level n. At v_k close to Ω/k, the exponential in (14) can be neglected and $\rho_m(v_x)$ is determined only by the factor $W(v_k)$. At velocities greatly differing from Ω/k,

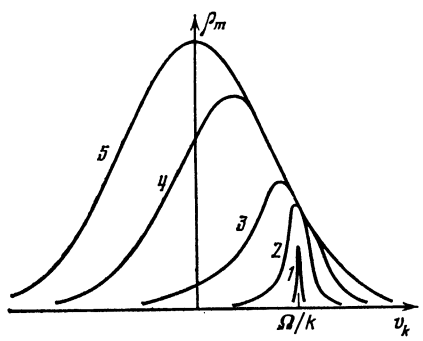

FIG. 1. Distribution function of the excited particles with respect to the velocity component along the wave vector under optical-pumping conditions, $(2\pi)^{1/2}\Gamma a/|v_x|=10^2$, $\Gamma/k\bar v=10^{-2}$, for different values of $\varkappa_{mn}^0$: 1—10^{-2}; 2—1; 3—10; 4—10^2; 5—∞.

$$|v_k-\Omega/k|\gg(\Gamma/k)\,(\varkappa_{mn}{}^0\Gamma a/|v_x|)^{1/2},$$

the argument of the exponential is small in absolute value and the optical pumping does not play a significant role, so that the Bennett peak has the usual Lorentz wings. For the width of the peak we have the estimate

$$\Delta v_k\sim(\Gamma/k)\,(\varkappa_{mn}{}^0\Gamma a/\bar v)^{1/2}. \tag{15}$$

We see that broadening is produced in this case there not only by the field but also by optical pumping.

To obtain numerical values of the current, we estimate the integral

$$\int Q(\mathbf{v})\,v_k\,d\mathbf{v}\sim\rho_m{}^2\sigma\bar v, \tag{16}$$

where ρ_m is the excited-particle density whose order of magnitude is $(N\bar v/\Gamma a)\,(\Delta v_k/\bar v)$. Substituting the expression for Δv_k and ρ_m in (16) and in (2), we obtain an estimate for the current:

$$J\sim eda\bar v\sigma\Gamma\varkappa_{mn}{}^0/v_s k^2. \tag{17}$$

Expressing ρ_m in (16) in terms of the absorbed power, we get

$$J\sim e\,(\Delta S/\hbar\omega\Gamma)^2\sigma\bar v^2/Vv_s, \tag{18}$$

which agrees with the estimate of Ref. 1 (V is the volume of the radiation-absorption region).

It is easy to estimate from (18) the value of the current under the typical experimental conditions $V=10^{-3}$ and $\Delta S=3$ mW. The published values of the cross section for reaction (3) range from $5\cdot10^{-18}$ to 10^{-15} cm^2 (Ref. 9). We use for the estimate $\sigma=10^{-16}$ cm^2. The order of magnitude of the light-induced current is 10^{-10} A.

As for the frequency dependence of the current, it can be stated that it is antisymmetric in Ω', and behaves at $\Omega'\gg k\bar v$ as the derivative of the squared absorption-line profile. In addition, since optical pumping broadens the Bennett structures, it is correct to conclude that the hyperfine structure is likewise not resolved in this case.

3. EXPERIMENT

We have investigated light-induced current in sodium vapor. The experimental setup consisted of a cw dye laser (Spectra-Physics model 375), a cell with sodium vapor, and a current meter.

The laser operated in a two-frequency regime, with a total lasing line width $\approx$480 MHz. The frequency was tuned by tilting an intracavity Fabry-Perot interferometer, comprising a plane-parallel quartz plate 3 mm thick. The typical radiation power was 70 mW. The radiation was focused by a lens into the working cell. The beam-constriction diameter was approximately 0.2 mm.

The working cell (Fig. 2) was a parallel-plate capacitor with openings at the electrode centers for passage of the beam. The distance between electrodes was 2.5 cm. They were housed in an evacuated glass container. Sapphire plates were placed between the glass windows and the electrodes. These plates, first, preserved the optical property of the windows during the experiment, since sapphire does not react

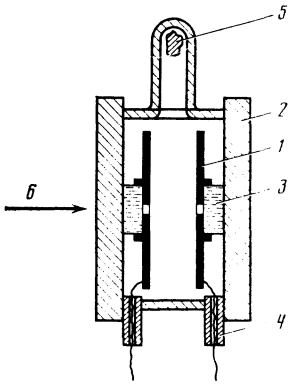

FIG. 2. Diagram of working cell: 1—electrodes; 2—glass windows; 3—sapphire plates; 4—beryllium-ceramic sleeves; 5—stub with metallic sodium, 6—laser beam.

chemically with sodium. Second, sapphire acted as an electric insulator. This was particularly important, since heating lowers the glass resistance appreciably, and furthermore an emf up to 0.5 V is produced on a glass-metal junction. This parasitic potential difference, greatly exceeding the ion energy, gathered the ions on one of the electrodes and made the signal symmetric about Ω'. This could be eliminated only by very carefully insulating the glass from the metal. The leads from the cell were insulated from the glass by beryllium ceramic sleeves.

The sodium vapor was fed to the cell from an autonomously heated stub. The cell parts were joined together with a heat-resistant adhesive. The vacuum was not worse than $2\cdot10^{-5}$ Torr. The cell working temperature was 130 °C. The stub was heated in a manner that produced a sodium-vapor density sufficient for noticeable absorption of the light. The vapor density was of the order of 10^{12} cm^{-3}, with $\approx$3 mW absorbed at the center of the line. The current signal was measured with a U5-6 dc amplifier.

The experiment consisted of recording the current and investigating its dependences on the radiation frequency, on the direction of the wave vector **k**. On the resonant-fluorescence intensity, on the voltage across the cell electrodes, and on the buffer-gas pressure.

A typical experimental plot of the current as a function of laser frequency is shown in Fig. 3. It can be seen that the plot is symmetric and there is no current at $\Omega'=0$. The sign

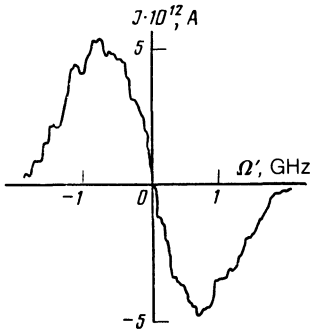

FIG. 3. Experimental plot of the current vs detuning of radiation frequency.

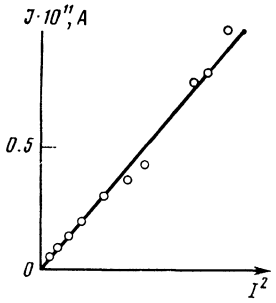

FIG. 4. Current vs squared intensity of the resonant fluorescence.

of the effect is such that the ion motion and the beam propagation are in the same direction for "blue" detuning of the frequency, and opposite for "red" detuning.

Reversal of the sign of the wave vector **k** changed also the direction of the current, likewise confirming our ideas concerning the nature of the effect. The light-induced current was of the order of 10^{-11} A, in agreement with the estimates of Sec. 2.

When a voltage was applied to the electrodes, the plot of the current as a function of Ω' became distorted. When the voltage reached $\approx$40 mV, the sign of the current no longer varied with the frequency detuning. This points to an ion energy $\approx$0.04 eV, i.e., of the order of the thermal energy of the atoms. With further increase of voltage, the current increased and saturated at $\approx$0.5 V. Saturation was reached when all the ions produced in the cell landed on one electrode, and all the electrons on the other. It is easy to calculate from the saturation current the total number of ions produced in the cell per unit time, and estimate the density of the charged particles. In our experiments the saturation current was $2\cdot10^{-10}$ A. The corresponding estimate for the ion density yields $\approx 4\cdot10^{-4}$ cm^{-3}. At this density the Debye radius is $\approx$1 cm. We operated therefore under near-plasma conditions.

Let us write the expression for the saturation current:

$$J_{\text{н}}=2ed\int Q(\mathbf{v})\,d\mathbf{v}. \tag{19}$$

Comparing (19) with (2), we easily find the relation between the saturation current and the light-induced current at equal $Q(\mathbf{v})$:

$$J=\frac{v_0}{4v_s}J_{\text{н}},\quad v_0=\int Q(\mathbf{v})v_k d\mathbf{v}\Big/\int Q(\mathbf{v})\,d\mathbf{v}. \tag{20}$$

Here v_0 has the meaning of the average initial ion velocity. From this equation and from the values of J and J_{sat} it follows that $v_0\approx v/2$. which agrees both with the estimates obtained from the value of the barrier voltage and with the theoretical value.

Evidence that the current is due to ions produced by associative ionization from the excited state is afforded by the quadratic dependence of the current on the intensity of the resonant fluorescence (see Fig. 4).

It can be seen from (2) that the light-induced current is rather simply related to the buffer-gas pressure. The effect can herefore serve as a rather simple method of measuring the frequency of the collisions of the ions with the buffer particles. Experimental plots of the current against helium and neon pressure are shown in Fig. 5. Our measured values of the diffusion coefficients of Na_2^+ ions in helium and neon at 400 K, when recalculated to 60 Torr, are respectively 0.53 ± 0.05 and 0.18 ± 0.04 cm^2/sec. The measurement results agree well with the handbook data. The diffusion coefficients of Na_2^+, calculated from the approximate equation of Ref. 10, are 0.55 and 0.19 cm^2/sec in helium and neon, respectively.

4. DISCUSSION AND CONCLUSIONS

We have reliably recorded the self-induced-current effect. The behavior of the effect agrees well with our ideas concerning its nature. We have qualitatively considered the behavior of an optically pumped three-level system in a collision-free situation. We note that the effect can be appreciably strengthened by using a two-frequency laser with a frequency difference equal to Δ, since this would obviate the need for optical pumping.

The effect can be use to study associative-ionization reactions, and to determine the cross sections and energy defects of these reactions. Using the velocity selectivity of the excitation, we can study the velocity dependence of the cross sections; this is of particular interest for endothermic reactions, in which the energy defect is compensated for by the translational-motion energy of the atoms. In additions, as shown in Ref. 1, light induced current can serve as a method of intra-Doppler spectroscopy for weak currents, in view of the accumulation of the slow ions. For the above reasons we regard further research into light-induced current as promising.

In conclusion, the authors thank S. P. Pod'yachev and D. A. Shapiro for valuable discussions during all stages of the work.

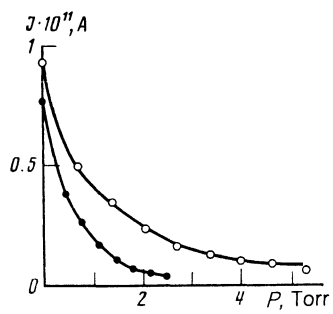

FIG. 5. Light-induced current vs helium (light circles) and neon (dark) pressure.

[1]S. N. Atutov, I. M. Ermolaev, and A. M. Shalagin, Pis'ma Zh. Eksp. Teor. Fiz. **40**, 374 (1984) [JETP Lett. **40**, 1185 (19843)].
[2]R. B. Green, R. A. Keller, G. G. Luther, *et al.*, Appl. Phys. Lett. **29**, 727 (1976)].
[3]F. Kh. Gel'mukhanov and A. M. Shalagin, Kvant. Elektron. (Moscow) **8**, 590 (1981) [Sov. J. Quant. Electron. **11**, 357 (1981)].

[4]A. I. Parkhomenko and V. E. Prokop'ev, Opt. Spektrosk. **53**, 1000 (1982).
[5]M. A. Vaksman and A. V. Gainer, Quant. Elektron. (Moscow) **9**, 901 (1982) [Sov. J. Quant. Electron. **12**, 570 (1982)].
[6]F. Kh. Gel'mukhanov and A. M. Shalagin, Pis'ma Zh. Eksp. Teor. Fiz. **29**, 773 (1979) [JETP Lett. **29**, 711 (1979)].
[7]I. M. Ermolaev, Avotmetriya No. 4, 80 (1985).
[8]V. R. Mironenko and A. M. Shalagin, Izv. AN SSSR Ser. Fiz. **45**, 995 (1981).
[9]A. N. Klyucharev and N. N. Bezuglov, Processes of Excitation and Ionization of Atoms by Light Absorption [in Russian], Leningrad Univ. Press, 1983, §4.2.
[10]Tables of Physical quantities, I. K. Kikoin, ed. [in Russian], Atomizdat, 1976, p. 291.

Translated by J. G. Adashko