populations of the low excited states influence the time evolution of the high-lying excited states of erbium even in a singly activated system. In fact, when the YOCl:Er^{3+} system is excited into the $^4S_{3/2}$ state (Fig. 3), multiphonon relaxation populates the $^4F_{9/2}$ level, from which luminescence with the proper lifetime begins. At the same time, all the remaining low-lying levels of erbium begin to be populated. This is due both the intracenter process and to cross relaxation processes. Therefore, after some time interval, the $^4F_{9/2}$ level starts to be populated not on account of relaxation from above, but on account of summation of the energies of two erbium ions, one of which is in the state $^4I_{11/2}$ and the other in the state $^4I_{13/2}$. The damping constant that will be observed at the later stages of the decay of the $^4F_{9/2}$ level is determined by the sum of the probabilities of the deactivation of the levels $^4I_{11/2}$ ($\tau = 2$ msec) and $^4I_{13/2}$ ($\tau = 5$ msec), yielding a value 1.4 msec. This is in good agreement with experiment.

[1)] It is appropriate to note that this result does not contradict the data of Ref. 6, inasmuch as under stationary excitation this relation does not hold and the second mechanism predominates.

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Relaxation currents in solids

V. K. Malinovskiĭ and B. I. Sturman

Institute of Automation and Electrometry, Siberian Division, USSR Academy of Sciences

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It is shown that in a crystal possessing a polar direction any process of relaxation to thermodyanmic equilibrium should be accompanied by a current. Unlike the currents connected with polarization direction ($\mathbf{j} = \dot{\mathbf{P}}$), relaxation currents do not destroy the existing charge separation but lead to the appearance and to maintenance of this separation. The nonuniformity that causes the charge separation can cause also electrization of the crystal for a prolonged time. Concrete models are considered, in which the disequilibrium of the crystal is attributed to an excess concentration of point defects. The question of activation of relaxation currents by light is considered.

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1. INTRODUCTION

The results of research into the photogalvanic effect in crystals that have no symmetry centers,[1,2] which has been intensively pursued in the past 4–5 years, point to the existence of a new class of transport phenomena. Namely, fluxes of charges, particles, and other physical quantities can exist in media without symmetry centers even in the absence of macroscopic inhomogeneities or external field. For these fluxes to exist it suffices only that there be no complete thermodynamic equilibrium in the medium. Thus, in the case of the photogalvanic effect a disequilibrium of the electron subsystem and a direct electric current are maintained by uniform illumination.

The disequilibrium of the medium need not necessarily be due to an external action. Even the presence of all types of defects (point, line, or volume) takes the medium quite far from thermodynamic equilibrium. The relaxation time of this disequilibrium obeys as a rule the activation law. Even at room temperatures, it can be quite long-on the order of months and years. Thus, a real crystal (as well as an amorphous medium) constitutes a highly unbalanced system that relaxes very slowly to equilibrium. It can be stated that in a medium that possesses at least one polar direction this relaxation should be accompanied by current flow. We emphasize forthwith that this effect differs substantially from the currents due to relaxation of an initially present charge separation (with polarization relaxation, $\mathbf{j}=\dot{\mathbf{P}}$ - the electret effect[3]). In contrast to the latter phenomenon, the relaxation process leads not to annihilation of the existing charge separation, but to the onset and maintenance of such a separation. The currents dealt with below are by no means time derivatives of the polarization.

We consider below simple concrete models that lead to currents due to relaxation of point defects in a crystal lattice vector having a polar vector **c**. We estimate the resultant electrostatic fields and discuss the question of activation of the relaxation process by light.

We present an *a priori* estimate of the relaxation current to gain an idea of the possible magnitude of the effect. Let the defect concentration, expressed in terms of point defects, be N, let the relaxation time be t_0, and let the directed displacement of the defect during the relaxation time be a ~10^{-7}–10^{-8} cm. Then

$$j \sim eNat_0^{-1}. \tag{1}$$

Putting $t_0 = 3\times10^7$ sec (one year), $N=10^{19}$ cm^{-3} (a prefectly realistic concentration in an ionic ferroelectric such as $LiNbO_3$), we get $j=10^{-13}$ A/cm^2. As applied to dielectrics, which are in fact dealt with here, such currents can give rise to rather large quasistationary electric fields, $E=j\sigma^{-1}$. Assuming a conductivity $\sigma=10^{-16}$ Ω/cm^2, we have $E\sim10^3$ V/cm.

It seems that mechanisms similar to those considered below can be realized not only in ferroelectric crystals, but also in amorphous media, and can be responsible for many electric and photoelectric phenomena.

2. BASIC MODEL

Consider a cubic crystal lattice containing arbitrarily distributed point defects, viz., vacancies and interstitial atoms, with respective concentrations N_V and N_I. The motion of such a lattice gas is a random diffusive process characterized by migration energies $\varepsilon_{V,I}$. The probability of a defect hopping over to a neighboring equivalent position is subject to the relation[4,5]

$$W_{V,I}=\Omega_{V,I}\exp(-\varepsilon_{V,I}/T); \tag{2}$$

$\Omega_{V,I}$ are of the order of the phonon frequencies. The diffusion coefficients $D_{V,I}$ for a cubic lattice are equal to

$$D_{V,I}={}^1/_6 W_{V,I}a^2,$$

where a is the lattice constant.

Under thermodynamic equilibrium, the defect concentration $N_{V,I}^T$ is determined by the known Frenkel' formulas.[5] We associate the disequilibrium of the medium with the excess defect concentration, i.e., we assume that $N_{V,I}\gg N_{V,I}^T$. Neglecting the interaction of the defects with the surface, we can assume that the mechanism of the relaxation to the equilibrium is annihilation of vacancy-atom pairs in the interstices, which leads to formation of defect-free cells. The inverse process, i.e., spontaneous production of defect pairs, can be neglected since $N_{V,I}^T$ is small. We assume for simplicity that the vacancies have the lowest migration energy, i.e., that they constitute the most mobile component. The interstitial atoms can therefore be regarded as immobile sinks, and one can refer to the vacancies as particles that wander over the lattice.

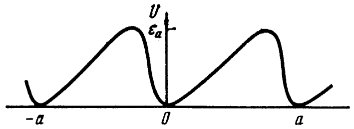

FIG. 1. Effective potential energy for particle.

So far, our assumption had nothing to do with the fact that the crystal has no central symmetry. The presence of a polar direction **c** causes the potential relief that determines the hopping of the particles to acquire a sawtooth shape (to become asymmetrical), see Fig. 1. Such an asymmetry, however, is still insufficient for a current to appear. The probabilities of hopping into right-hand and left-hand equivalent positions are equal. Otherwise we would obviously have current under conditions of total thermodynamic equilibrium.

We examine now the effect of the asymmetry of the potential on the annihilation of a pair of defects. In a crystal having a polar direction, the vacancy positions on the right and on the left of the interstitial atoms are equivalent. Accordingly, the potential barriers that must be overcome for the recombination from the right and left positions should be different, see Fig. 2.

Calculation of the activation energies is a complicated problem of many-body theory.[4] It can be simplified by assuming that the difference between the activation energies is due to the difference between the dipole moments of the cells with and without defects. The additional dipole moment facilitates hopping from the left and hinders it from the right (Fig. 2). If the activation-energy difference $\Delta\varepsilon > T$, then pair annihilation proceeds essentially asymmetrically.[1] We assume for the sake of argument that the largest hopping probability W_+ is possessed by an ion on the left of the sink, i.e.,

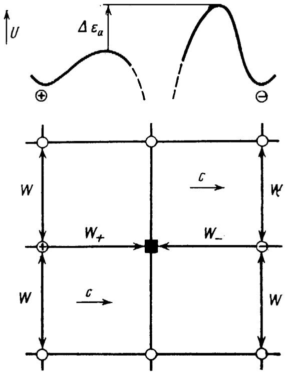

FIG. 2. Illustration of migration and annihilation of defects in a cubic lattice. The square marks an interstitial atom (sink).

the annihilation act is accompanied by a hop in the c direction. The annihilation is followed by formation of a regular unit cell, and the ions that make it up no longer migrate.

The model formulated does not use in essence any particular assumptions concerning the properties of the crystal. It merely reflects the fact that the crystal has no central symmetry.

We arrive thus at the following picture of the carrier motion: The ions wander randomly over the lattice sites until the ions land on sites to the left of the sink. This is followed by either irreversible annihilation or by continuation of the wandering.

We write down now the kinetic equations for our model. The rate of change of particle concentration is

$$dN/dt = -W_+ n_+, \tag{3}$$

where n_+ is the number of particles per unit volume on the sites + (Fig. 2). We must recognize next that the probability of finding the particle on some particular site depends generally speaking on the site position relative to the capture centers. In the absence of long-range action, however, with increasing distance from the nearest center the probability should rapidly tend to the mean value $N_I N_0^{-1}$, where $N_0 \equiv a^{-3}$ is the total number of sites per unit volume. If the number z of the nearest neighbors is large (the coordination number for a cubic lattice is $z=6$) then, accurate to z^{-1}, the population of all sites with the exception of the + sites can be regarded as equal to $\bar{n} = N_I N_V / N_0$. We can therefore write

$$\frac{dn_+}{dt} = -(W_+ + 5W)n_+ + 5WN_V\frac{(N_V+\Delta)}{N_0}. \tag{4}$$

We have taken into account in (4) the fact that the defect-concentration difference $\Delta = N_I - N_V$ is conserved in the course of relaxation in our model.

The system (3) and (4) requires, besides knowledge of the initial defect concentration, also of the initial condition for n_+. It can be stated that by specifying $n_+(t=0)$ we take into account the correlation of the initial particle distribution, a correlation that reflects the prior history of the crystal. Actually, as seen from (4), within a time $\tau \sim (W_+ + 5W)^{-1}$ the quantity n_+ relaxes to the value

$$n_+ = 5WN_V(N_V+\Delta)/N_0(W_+ + 5W). \tag{5}$$

If $W_+ \gg 5W$, then $n_+ \ll \bar{n}$; in the opposite case $n_+ \approx \bar{n}$.

The characteristic time t_0 of the relaxation of the total particle concentration will be shown below to be much longer than τ. Being interested in just such long times $t \gg \tau$, i.e., disregarding the short initial stage of the relaxation, we have from (3)–(5)

$$\frac{dN_V}{dt} = -\nu\frac{N_V(N_V+\Delta)}{N_0}, \quad \nu = \frac{5WW_+}{W_+ + 5W}. \tag{6}$$

These solutions are obviously meaningful in the region $N_V, N_V + \Delta > 0$.

It follows from (6) that the relaxation kinetics depends substantially both on the ratio of the hopping probabilities W and W_+, and on the ratio of the concentrations N_I and N_V. A characteristic feature of the solutions (6) is the dependence of the relaxation time on the defect concentration. If $N_I \gg N_V$, then N_V decreases exponentially with a time $t_0 = N_0/\nu\Delta$. If $\Delta = 0$, then

$$\frac{N_V}{N_V(0)} = \left(1 + \nu\frac{N_V(0)}{N_0}t\right)^{-1}. \tag{7}$$

If $N_V \gg N_I$, then N_V decreases exponentially to the level $|\Delta|$ with a time $t_0 = N_0/N_V(0)\nu$. The quantity ν, as expected, is determined by the smaller of the probabilities W_+ and $5W$, i.e., by the slowest process that leads to relaxation.

We determine now the current that accompanies the relaxation to equilibrium. Since the only contribution to the current is made by particles the execute hops from the sites ±,

$$\mathbf{j} = ea(\bar{n}W + n_+ W_+)\mathbf{c}. \tag{8}$$

Introducing for convenience the relative defect concentration $C = N_V/N_0$ and using (5), we obtain

$$j = \frac{eW}{a^2}\eta C\left(C + \frac{\Delta}{N_0}\right), \tag{9}$$

where $\eta = (5W + 6W_+)(5W + W_+)^{-1}$ is a parameter that ranges from 1 to 6.

We estimate now the current (9). Putting $\Delta = 0$ and $W_+ \gg W$, we have

$$j \approx 6Ce/a^2 t_0, \quad t_0 = (WC)^{-1}, \tag{10}$$

which obviously agrees with (1). Putting $C = 10^{-3}$, $a = 10^{-8}$ cm and $t_0 = 10^7$ sec we get

$$j \sim 10^{-12}\ \text{A/cm}^2. \tag{11}$$

The migration energy ε_V corresponding to these data is approximately 1.2 eV.

The law governing the decrease of the current is described by Eqs. (6) and (7). The slowest decrease at $\Delta = 0$, and then $j \propto (1 + t/t_0)^{-2}$. In the opposite case, the decrease is exponential.

Let us determine the maximum electrostatic field possible in the considered model. If the Maxwellian time $t_M \equiv \varepsilon_0/4\pi\sigma \ll t_0$, then a quasistationary field $\mathbf{E} = -\mathbf{j}\sigma^{-1}$ is established inside an isolated crystal. Using the Einstein relation for the mobility, $D = \mu T/e$, we get

$$\sigma = e^2 W a^2 N_V / 6T. \tag{12}$$

The condition $t_M/t_0 \ll 1$ necessary to establish a quasistationary field can be written, as follows from (10) and (12), in the form $\varepsilon_0 aT/e^2 \ll 1$. This condition is almost always satisfied. From (10) and (12) we have for the maximum field

$$E = 6CT/ea. \tag{13}$$

At room temperature, according to (13), we have $E = 2 \times 10^7 C$ [W/cm]. Thus, owing to the low (ionic) conductivity, small relaxation currents are capable of producing appreciable fields. As follows from (13), the field E increases with rising temperature. With rising T, however, the lifetime of the nonequilibrium state decreases drastically.

We emphasize that in the considered model the role of the interstices as asymmetrical sinks for the vacancies is in no way singled out. Any other defect can serve as an asymmetrical sink, e.g., impurity atom, various complexes, or violation of the stoichiometry of the composition. It can be stated that any reaction similar to that considered above[2] should be accompanied by current flow if the crystal matrix in which the reaction takes place has a singled-out direction.

It must be noted that the presence of relaxation currents leads also to modification of the phenomenological equations that describe the migration processes. In addition to diffusion of particles and their recombination, the phenomenological equations should contain terms that describe the drift:

$$\frac{\partial N}{\partial t}=\left(\frac{\partial N}{\partial t}\right)_{\mathrm{rec}}-\operatorname{div}\frac{\mathbf{j}}{e}+D\Delta N. \tag{14}$$

The additional term $\operatorname{div}(\mathbf{j}/e)$ takes into account the existence of a flux of particles in the absence of a concentration gradient.

3. ELECTRONIC-CURRENT MECHANISMS

We have considered above the relaxation current due to motion of ions. Mechanisms are possible, however, in which the charge carriers are electrons, and the ion system serves only as the source of the disequilibrium.

Point defects (vacancies, impurities, and others) can cause deep energy levels in the forbidden band. In the course of annihilation or coalescence of defects, the energy spectrum becomes substantially restructured. In particular, the localized electronic states can vanish partially or completely. An electron in a strongly localized state prior to the recombination act becomes less localized, free, or else captured by close-lying centers. We consider the limiting case, when the localized state vanishes completely after the recombination of a pair of defects. The change of the state of the electron in such a process can be described by a potential energy $U^e(\mathbf{r},t)$ that decreases slowly in time. The characteristic time τ_0 at which the potential vanishes is determined by the motions of the ions; it can be assumed that $\tau_0\sim\Omega^{-1}$ [see (2)].

In media without a symmetry center, the potential $U^e(\mathbf{r},t)$ is obviously likewise noncentral, see Fig. 3. So long as the populated energy level is deep enough in the potential well.

$$E(t)\gg\hbar\Omega,$$

the change of the potential can be regarded as adiabatic.[9] During this longest stage, the electron-localization radius increases to a value

$$l\sim(\hbar/2m\Omega)^{1/2}\gg a. \tag{15}$$

By virtue of the asymmetry of the potential, such a delocalization shifts the center of gravity of the electron wave packet to the right (Fig. 3). If there is no reason for assuming the potential hump to be small, then the size of this shift is of the order of l.

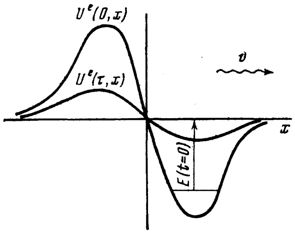

FIG. 3. Destruction of localized state and asymmetric ionization of electrons following coalescence of defects.

Next, within a time $\Delta t\ll\tau_0$, the decrease of the potential is no longer adiabatic and is accompanied by a total vanishing of the discrete spectrum. The result is a packet with a momentum width $\Delta k\sim 2\pi l^{-1}$. Waves with $k<0$ will be scattered by the residue of the potential,[3] and consequently the average momentum of the electrons will differ from zero at $t\gg\tau$. The foregoing reasoning holds also for the three-dimensional case.

The electron produced after the recombination act has thus a characteristic directional velocity

$$\mathbf{v}=\xi(\hbar\Omega/2m)^{1/2}\mathbf{c}. \tag{16}$$

The asymmetry parameter $\xi\lesssim 1$ depends on the details of the structure of the potential during the concluding stage of its vanishing.

It is easy to estimate the total current in the considered model. The current has two components corresponding to the shift of the gravity centers of the localized states and to the motion of the free electrons. Denoting the electron momentum isotropization time by Γ^{-1}, we have

$$j\approx e\left|\frac{dN}{dt}\right|\left[\left(\frac{\hbar}{2m\Omega}\right)^{1/2}+\frac{\xi}{\Gamma}\left(\frac{\hbar\Omega}{2m}\right)^{1/2}\right]. \tag{17}$$

Comparing (17) with (6) and (9), we see that other conditions being equal the electron current exceeds the ion current if

$$l_{\mathrm{eff}}=l+\xi v/\Gamma>a.$$

Putting $\Omega=10^{14}\ \mathrm{sec}^{-1}$, $\xi=10^{-1}$, and $\Gamma=10^{13}\ \mathrm{sec}^{-1}$, we obtain $l_{\mathrm{eff}}=3\times10^{-7}$ cm, i.e., the current increases by more than one order of magnitude compared with (11).

The foregoing estimate of the current presupposes that $l\ll R$, where R is the mean distance to the nearest capture center. If $l\gtrsim R$, then the electron can tunnel to this center and no free electron appears; in this case

$$j\gtrsim eR\,|dN/dt|.$$

We emphasize that in the considered model the total charge

$$Q=\int_0 j\,dt,$$

flowing through the crystal surface is, generally speaking, not a characteristic of the ion subsystem; Q can be arbitrary to a considerable degree.

4. ACTIVATION OF THE RELAXATION CURRENTS BY LIGHT

Illumination of the crystal can increase appreciably the current, by shortening the relaxation time τ. It can be stated figuratively that the light acts as a catalyst for the relaxation process. As seen from (2) and (9), the relaxation currents are most sensitive to changes of the temperature and of the activation energy.

Let us discuss the possibility of changing these quantities by applying light. The heating of the crystal by the light is as a rule small and of no interest. It must be borne in mind, however, that when light is absorbed by local centers the crystal heating can proceed in two stages. The electrons excited by the light to low lying levels or into the conduction band can release their energy in the immediate vicinity of the defect, via nonradiative recombination to a local level, and thus produce local heating.[10,11] The local heating can be due to the presence of local or pseudolocal phonon modes in the crystal with defects. The time of relaxation of this perturbation, τ_{loc}, can apparently exceed by several orders of magnitude the reciprical phonon frequencies,[10,12] If the electron recombination time $\tau_e \lesssim \tau_{loc}$, and the number of atoms effectively participating in the local or nonlocal mode is n_{ph} ($n_{ph} \gg 1$), then the heat rise is

$$\Delta T \approx \hbar\omega / n_{ph}. \tag{18}$$

Under stationary conditions the number of overheated centers is

$$\delta N \approx \frac{\varkappa J}{\hbar\omega} \tau_{loc},$$

where $\varkappa$ is the absorption coefficient, and J is the light intensity. Consequently, the change of the current is

$$\delta j = G^{\tau} j, \quad G^{\tau} = \frac{\varkappa J}{\hbar\omega} \frac{\tau_{loc}}{N} \exp\left(\frac{\hbar\omega\varepsilon_a}{T^2 n_{ph}}\right). \tag{19}$$

The most important here is the dependence on the parameter n_{ph}. Assuming the estimates $\varkappa = 1$ cm^{-1}, $J = 1$ W/cm^2, $\tau_{loc} = 10^{-11}$ sec, $T = 300$ K, and $\hbar\omega = \varepsilon_a = 1.5$ eV, we obtain by way of the condition $\delta j \gg j$

$$n_{ph} < 60. \tag{20}$$

This estimate is admittedly quite vague, thus indicating that the considered mechanism is realistic. We call attention also to the linear dependence of the current δj on the light intensity and to its sensitivity to changes of temperature.

The other way of increasing the current is to decrease the migration energy by means of light. Since ε_a depends essentially on the electronic state of the defect, a state that determines the chemical bonds, the use of light to excite the electron to higher levels can change significantly the activation energy (up to inversion of the potential barrier[13]). The concrete realization of these changes can vary: the excited electrons can become localized near saddle points, cause charge exchange of diffusing atoms, or perturb the chemical bonds (see Refs. 11 and 13 and the literature cited therein). Let the decrease of the migration energy be τ_l, and the lifetime of the excited state $\Delta\varepsilon_a$. Then, in analogy with the reasoning above, we get

$$\delta j = G^{a} j, \quad G^{a} = \frac{\varkappa J}{\hbar\omega} \frac{\tau_l}{N} \exp\left(\frac{\Delta\varepsilon_a}{T}\right). \tag{21}$$

For the parameters assumed above and $\tau_l = 10^{-9}$ sec, we obtain in lieu of (20)

$$\Delta\varepsilon_a > 0.5 \text{eV}.$$

5. CONCLUSION

The models considered above can, of course, not claim to permit profound and complete research into relaxation currents. They represent in essence a concrete form of the general statement that currents can flow in a nonequilibrium medium with a polar direction even in the absence of external action or spatial inhomogeneity.

It is clear also that the considered asymmetric recombination models are not peculiar to point defects only. They can be extended without significant changes to more complex processes such as migration of point defects to dislocations, to grain boundaries, to domain walls, etc. Mechanisms related to those considered here can apparently be realized not only in crystalline but also in amorphous media without symmetry centers, in which the relaxation process can be of the type of structural ordering, crystallization, or chemical transformation. More extremal characteristics of the relaxation currents can be expected in such cases.

We note also that the signs of the relaxation currents are as a rule not definite. In principle, different recombination channels can make contribution in both the c and –c directions.

We consider now the question of observing the relaxation currents. This question has several aspects. The effect of preservation of the internal electric field for a long time has been known for a large number of substances[3] (both amorphous and crystalline). This phenomenon was previously attributed to relaxation of initially effected polarization (to the electret effect). The possibility is not excluded that in many cases this phenomenon can be attributed to relaxation currents.

In the case of ferroelectrics of the $LiNbO_3$ type the relaxation currents might explain in principle the existence of Chen fields, which are responsible, together with the photogalvanic effect, for the photorefraction phenomenon.[14]

Finally, one more possibility can be suggested. It is known that constant photocurrents (lasting for several hours) at a level $\sim 10^{-10}$ A/cm^2 at $J \sim 0.1$ W/cm^2 are presently observed in many pyroelectrics[1,2] and are interpreted as photogalvanic currents. From the theoretical point of view, it is very strange that the observed currents do not have to a full extent those polarization properties that follow from the crystal symmetry. As a rule, the photogalvanic tensor β_{ijl} in pyroelectrics only one independent component $\beta_{ijl} = \beta c_i \delta_{jl}$ corresponding to current along the polar axis of the crystal. Yet in semiconducting non-ferroelectric crystals there is a good correspondence between the crystal symmetry and the properties of the photocurrents.[2] As already noted, activated relaxation currents are linear in the light intensity. If relaxation currents $j \sim 10^{-12}$ A/cm^2 can exist for a year, then activated currents $j \sim 10^{-10}$ A/cm^2 can be preserved almost constant for many hours. Although on the whole it is apparently impossible to attribute the experimental data on the photogalvanic effect to relaxation currents, the question of their role in observable photoelectric phenomena remains open.

1) We assume in fact short-range forces of interaction between the defects. It is this which justifies the use of the lattice model (see also Refs. 6–8). Coulomb attraction can only accelerate the relaxation to equilibrium.

2) The relaxation process considered above can be regarded as a modification of a diffusion-limited reaction in the solid phase.[6]

3) It is easy to verify that by the instant of vanishing of the discrete spectrum the coefficient of reflection from the barrier is generally speaking not yet small.

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The drag of electron-hole drops by ultrasound

A. S. Alekseev, T. I. Gal'kina, V. N. Maslennikov, and S. G. Tikhodeev

P. N. Lebedev Physical Institute, Academy of Sciences USSR

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The drag of electron-hole drops (EHD) in germanium over macroscopic distances up to 4 mm, induced by traveling ultrasonic waves, is observed. The drag effect is detected by using a $\lambda = 3.39\ \mu$m probing laser beam. The absorption change in the probing region, which is proportional to the change in the carrier density in the region, is measured. A theoretical model is developed to describe EHD drag by ultrasound, in which the generation of the carriers and their finite lifetimes are taken into account. The model permits a semiquantitative explanation of the dependence of the effect on the intensity of the ultrasonic pulse and also on its duration and on the temperature during the experiment.

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The electron-hole liquid in pure undeformed germanium is formed in the shape of a cloud of fine (with radius 10^{-4}–10^{-3} cm) electron-hole drops (EHD).[1] Another possible method of motion of the drops is the action on them of ultrasound.[5,6] We have previously observed the change in the ultrasonic attentuation in photo-excited germanium with increase in the sound intensity[7] and the "flareup" of recombination radiation of the EHD upon its interaction with the ultrasound.[8] These effects have been explained with the help of representations on the EHD drag by the ultrasonic wave.

The basic purpose of the present research was to obtain direct experimental proof of the EHD drag by the ultrasound predicted in Ref. 5. Theoretically, the drag of the drops of the electron-hole liquid by the ultrasound is considered in the first section of the present paper. In the second section, we describe the experimental method, and in the third, the experimental results that were obtained are set forth and discussed.

I. DRAG OF THE EHD BY ULTRASOUND. REDISTRIBUTION OF THE EHD CLOUD DENSITY

Let current carriers with velocity g be generated continuously at some region of the semiconductor. We shall assume that they are all quickly condensed (within 5×10^{-7} sec[9]) into EHD, which form a cloud in the region of generation with drop density N and mean density of the condensed carriers in the cloud $n=\frac{4}{3}\pi R^3 N n_0$, where n_0 is the density of the liquid in the EHD, R is the radius of the EHD. We now consider how the mean density of the cloud of drops n changes under the action of a longitudinal ultrasonic wave of frequency ω, with a wave vector $\mathbf{k}$ directed along the x axis, and a wavelength that is large in comparison with the dimensions of the drops $\lambda \gg R$, $\lambda = 2\pi/|\mathbf{k}|$. In the deformation field of such a wave, $\varepsilon=\varepsilon_0\sin(\omega t-\mathbf{kr})$, where $\varepsilon=\operatorname{div}\mathbf{u}$, $\mathbf{u}(\mathbf{r},t)$ is the displacement of the point $\mathbf{r}$ of the semiconductor from the equilibrium position, the force[5]