preference to an explanation of the PICMP with the model using "near" and "far" sites with Fe^{2+} ions,[2,9] a model that takes account also of the stabilization of the DW.

Our results of the investigation of the photoinduced stabilization of the DW agree with the results of Halsma and Robertson,[16] who observed directly the effect of light on the DW mobility and have shown that the DW light stops the motion of the DW in weak alternating magnetic fields. In contrast to the previously observed PICMP, which proceed smoothly without jumps and in which the light influences the properties (mobility) of the DW but does not change the domain structure, the photoinduced jump observed by us should cause also a change in the domain structure of the crystal.

The photoinduced increase of the higher harmonics of the total magnetic permeability, observed here for the first time, cannot be regarded as an obvious consequence of the previously observed photoinduced change in the rectangularity coefficient of the hysteresis loop. There exists a region of small values of H_m in which light decreases both the fundamental and the higher harmonics of μ_t. The photoinduced increase of the contribution of the higher harmonics to the total permeability attests to light-induced distortion of the symmetry of the potential wells and can be taken into account by introducing nonlinear terms in the equation of motion of the DW.

We note that the role of the DW pinning centers can probably be played by magnetic polarons, the possible formation of which under the influence of light was indicated by Belov, Koroleva, and Batorova.[17]

The results are of importance for a more complete understanding of the mechanism of the PICMP and its practical applications. The observed phenomena can be used to investigate the dynamics of domain walls and to analyze the crystal defects that serve as DW pinning centers.

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Nature of the dislocation charge in ZnSe

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A physical model is proposed to explain the experimentally observed anomalously large electric charges of moving dislocations in II–VI semiconductors. The model is based on the idea that broken bonds in the core of a dislocation are filled with electrons from point centers swept through by the dislocation during its motion. The theoretical predictions are compared with the experimental data for ZnSe.

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INTRODUCTION

The presence of electric charges at dislocations has been detected experimentally in many II-VI compounds: ZnS,[1-4] ZnSe,[4] CdS,[5] and CdSe.[4] A surprising feature is the very high linear density q of such charges, reaching one electronic charge per interatomic distance. The following interesting physical phenomena are associated with the motion of such strongly charged dislocations: the photoplastic effect,[6,7] deformation-induced luminescence,[2,8] influence of electrical boundary conditions on plastic deformation processes,[1] electroplastic effect,[9,10] and influence of dislocation motion on conduction current and photocurrent.[10,11] However, in spite of the importance of the nature of such high dislocation charges in these physical phenomena, the magnitude of the charge is accepted—with some exceptions[3,5]—in the cited papers simply as an experimental fact without interpretation. Osip'yan and Petrenko[3] used the results of an experimental investigation of the influence of illumination of the purity of ZnSe single crystals on the magnitude of the dislocation charge q to put forward the hypothesis that q might be affected by the exchange of electrons between a dislocation and point defects swept through by the dislocation during its motion.

Zaretskiĭ *et al.*[5] investigated the charges carried by dislocations of various types in CdS and demonstrated that only dislocations with broken bonds in their cores are charged. The present paper is devoted entirely to the nature of dislocation charges in II-VI compounds. In the first part of the paper, we shall put forward a physical model of the formation of dislocation charges and, in the second part, we shall compare the predictions deduced from this model with the experimental results.

PHYSICAL MODEL

It follows from the published results[1,4,5] that, among dislocations in II-VI compounds, only those are charged which have broken bonds in their cores. Therefore, the presence of such bonds will be regarded as the cause of the appearance of electric charges. As demonstrated by Read,[12,13] the presence of broken bonds in a dislocation core results in the capture of electrons by these bonds, produces a dislocation energy level E_d in the forbidden band of a semiconductor, and generates electric charges along the dislocation line. The validity of these ideas is now supported by an enormous amount of experimental material obtained in investigations of elemental material obtained in investigations of elemental semiconductors (Ge, Si) and of compounds (GaAs, InSb, etc.). The electron occupancy of a dislocation level is governed by the equilibrium Fermi distribution

$$f=\left(1+\exp\left(\frac{E_d(f)-\mu}{T}\right)\right)^{-1}, \tag{1}$$

where $f=a/l$ is the occupancy, a is the distance between the broken bonds (equal to the lattice constant), and l is the distance between the electrons captured by the bonds.

The Coulomb repulsion between electrons in the dislocation level E_d causes this level to rise with the electron occupancy so that it very rapidly reaches the chemical potential μ and then captures no further electrons. This behavior is responsible for the small experimental values of the occupancy $f \lesssim 0.1$ of dislocations at rest.

However, in the case of a moving dislocation, we cannot calculate f using the equilibrium distribution function (1). One of the reasons for this is—as will be shown in Appendix II—that the time is too short for the electrons. localized at point defects to reach thermal equilibrium with the conduction band when a charged dislocation passes nearby. Consequently, the occupancy of energy levels near a dislocation (including the level E_d) is no longer governed by the thermal equilibrium conditions and we can find the distribution of electrons between the various energy levels only by solving the kinetic equations describing electron exchange between them.

In the coordinate system (Fig. 1) linked to a dislocation moving at a constant velocity v_d these equations are independent of time:

$$P_{dc}-P_{vd}=\sum_t \iint_{-\infty} dx\,dy\left\{\omega_{td}n_t(x,y)(1-f)\psi(x,y)\Big/\left[1+\exp\left(\frac{E_d-E_t(x,y)}{T}\right)\right]\right.$$
$$\left.-\omega_{td}(N_t-n_t(x,y))f\psi(x,y)\Big/\left[1+\exp\left(\frac{E_t(x,y)-E_d}{T}\right)\right]\right\}, \tag{2}$$

$$\frac{dn_t(x,y)}{dx}=\frac{1}{v_d}\left\{-n_t(x,y)P_{tc}-\omega_{td}n_t(xy)(1-f)\psi(x,y)\Big/\left[1+\exp\left(\frac{E_d-E_t(x,y)}{T}\right)\right]\right.$$
$$\left.+\omega_{td}(N_t-n_t(x,y))f\psi(x,y)\Big/\left[1+\exp\left(\frac{E_t(x,y)-E_d}{T}\right)\right]\right\}, \tag{3}$$

where x is the coordinate along the direction of motion of the dislocation; y is the impact parameter (Fig. 1); ω_{td} is the attempt frequency of transitions between a local center of energy $E_t(x,y)$ and a dislocation level; N_t and $n_t(x,y)$ are, respectively, the concentration of such centers and the density of electrons at these centers; $\psi(x,y)$ represents the tunneling of electrons from a point defect to a dislocation or back again and—to within the precision required—is of the form $\exp\{-(x^2+y^2)^{1/2}/r_{0t}\}$; since in high-resistivity semiconductors only the deep strongly localized states are filled with electrons, it is found that r_{0t} is of the order of the interatomic distance; P_{dc} and P_{vd} are the electron fluxes from a dislocation to the conduction band and from the valence band to a dislocation; P_{tc} is the probability of release (per unit time) of an electron from an E_t center to the conduction band. The summation in Eq. (2) is carried out over the various types of impurity center.

The first term in the integrand in Eq. (2) describes the flux of electrons from a center E_t at a point (x,y) to a dislocation; the second term represents the flux in the opposite direction. The flux of electrons from a unit length of a dislocation to the conduction band P_{dc} was calculated by Shikin and Shikina[14]:

$$P_{dc}=f\omega_{dc}\left(\frac{\pi e^2 f}{T\varepsilon a^3}\right)^{1/2}\exp\left(\frac{E_d+Af(\ln(a'f^{3/2}/T)-1)}{T}\right\}, \tag{4}$$

$$A=\frac{2e^2}{\varepsilon a},\quad a'=\frac{e\hbar}{(\varepsilon m a^3\pi)^{1/2}},\quad \omega_{dc}\sim\frac{E_c-E_d}{\hbar}. \tag{5}$$

If

$$E_G-2|E_d|+Af(\ln(a'f^{3/2}/T)-1)\gg T \tag{6}$$

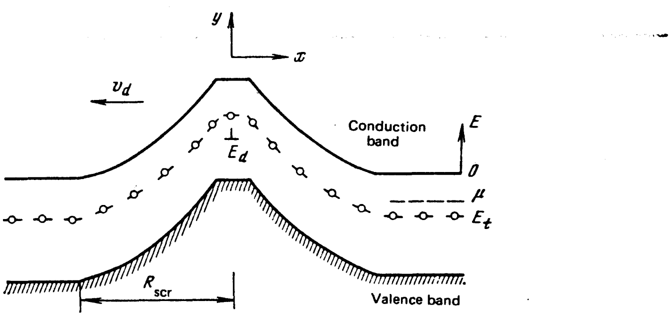

FIG. 1. Schematic representation of the bending of the allowed energy bands near a dislocation and the coordinate system adopted in calculations.

(E_G is the forbidden band width), the flux P_{vd} in Eq. (2) can be ignored compared with P_{dc}. We shall show in Appendix II that, for all the electron-filled E_t centers ($E_t \lesssim \mu$), we can also ignore the flux P_{tc} in Eq. (3). In deriving Eqs. (2) and (3), we have ignored electron transitions between the E_t centers because they are separated by large distances ($\gtrsim$100 Å). Finally, we have ignored electron transitions from the conduction band to point defects and dislocations because, in high-resistivity semiconductors (only in such semiconductors can we determine f from the dislocation current[3]), the conduction-electron density is very low ($\lesssim 10^6$ cm^{-3}) even far from a dislocation, whereas near a dislocation the bending (rise) of the bands by an amount $\sim Af\ln(R_{scr}/a) \gg T$ makes this density completely negligible.

Subject to these comments, the system (2)–(3) becomes

$$\sum_t \iint_{-\infty}^{+\infty} dx\, dy\, g_t(x,y)\left\{n_t(x,y)(1-f)-(N_t-n_t(x,y))f\exp\left(\frac{E_d-E_t(x,y)}{T}\right)\right\}$$
$$=\omega_{dc}\left(\frac{\pi e^2 f^3}{T\varepsilon a^3}\right)^{1/2}\exp\left\{\frac{E_d+Af(\ln(a'f^{1/2}/T)-1)}{T}\right\}, \tag{7}$$

$$\frac{dn_t(x,y)}{dx}=\frac{g_t(x,y)}{v_d}\left\{-n_t(x,y)(1-f)+(N_t-n_t(x,y))f\exp\left(\frac{E_d-E_t(x,y)}{T}\right)\right\} \tag{8}$$

where

$$g_t(x,y)=\omega_{td}\exp\left[-\frac{(x^2+y^2)^{1/2}}{r_{0t}}\right]\Big/\left\{1+\exp\left[\frac{E_d-E_t(x,y)}{T}\right]\right\}. \tag{9}$$

The solution of the system of the steady-state equations (7) and (8) is given in Appendix I:

$$(1-f)v_d\left(\sum_t{}' n_t(-\infty)R_{1t}\right)=\omega_{dc}\left(\frac{\pi e^2 f^3}{T\varepsilon a^3}\right)^{1/2}\exp\left\{\frac{E_d+Af[\ln(a'f^{1/2}/T)-1]}{T}\right\}, \tag{10}$$

where the summation is carried out over all the levels E_t satisfy the condition

$$E_t-E_d \geqslant Af\left[\ln\left(\frac{r_{0t}}{a}f\ln\left(\frac{\omega_{td}r_{0t}}{v_d}\right)\right)\right]\approx Af\approx 0.4\ \text{eV} \tag{11}$$

(for $f\approx 0.5$) and $R_{1t}\sim 10^{-6}$ cm.

We shall rewrite Eq. (10) in the form

$$f=\frac{1}{A[\ln(a'f^{1/2}/T)-1]}\left[-E_d+T\ln\left\{\frac{v_d(1-f)}{\omega_{dc}(\pi e^2 f^3/T\varepsilon a^3)^{1/2}}\sum{}' n_t R_{1t}\right\}\right]. \tag{12}$$

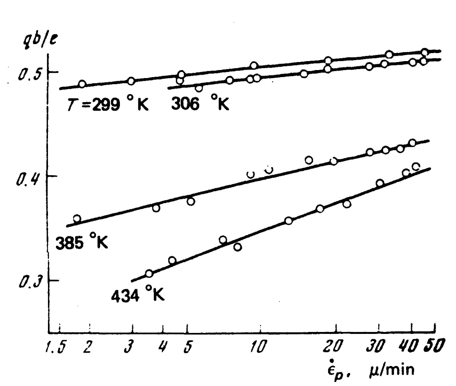

FIG. 2. Dependences of the dislocation charge $qb/e = f$ in ZnSe on the deformation rate.

Since the expressions in the logarithm of Eq. (12) are much larger than unity, the main dependences of f on v_d and T will be practically linear when plotted against $\ln v_d$ (T = const) and against T (v_d = const). Therefore, in the experimental check on the theory, we shall concentrate our attention on the velocity and temperature dependences of f. Moreover, on increase in $\sum' n_t$ by several orders of magnitude (which occurs as a result of illumination with interband light), we may expect a considerable increase in f. Finally, in the experimental part, we shall compare the occupancy f, deduced from the equilibrium model in Read [Eq. (1)], with the value of f deduced from Eq. (12) for parameters typical of ZnSe.

METHOD

We used n-type melt-grown ZnSe with the sphalerite structure and a room-temperature resistivity $\rho = 10^8 - 10^{13}$ Ω cm. Samples of 6 × 4 × 1.5 mm were cut from ingots by a diamond saw and were then ground with abrasive powders and polished with a diamond paste; next, they were subjected to chemical polishing in a solution of CrO_3 in HCl to remove the cold-worked layer. Indium contacts were deposited on the 6 × 4 mm face and the ohmic nature of these contacts was specially checked. The 1.5 × 4 mm face was parallel to the ($0\bar{1}1$) plane and the (111) plane made an angle of 45° with the long edge (6 mm) of the sample.

A sample was deformed by compression at a constant rate parallel to the long edge. The deformation rate was varied from 0.5 to 1000 μ/min. Selective etching in a mixture of 6 g NaOH + 4 g $K_3[Fe(CN)_6]$ + 50 g H_2O showed that plastic derormation was due to the motion of dislocations in a single plane (111), coinciding with the plane of stacking faults and making an angle of 45° with the long edge of the sample. A study of the dislocations was at least two orders of magnitude higher than that of the β locations, so that the measured charges were identical, to within ~1%, with the charges of the α dislocations.

The dislocation charges were determined by the method of dislocation currents from the ratio of such a current to the rate of plastic deformation $\dot\varepsilon_p$. This method was described in detail elsewhere.[1,3,5] We selected those ZnSe ingots for which, beginning from deformations of 2–3%, the value of f was independent of the degree of deformation (up to ≈15%).

EXPERIMENTAL RESULTS

Figure 2 shows the dependences of the dislocation charge in ZnSe at various temperatures on the plastic deformation rate $\dot\varepsilon_p \propto v_d$. The dislocation charge is reduced to the electronic charge e and multiplied by the distance between the broken bonds b, so that $qb/e = f$. The $\dot\varepsilon_p$ axis is logarithmic. In agreement with Eq. (10), we found that, in terms of these coordinates, f was a linear function of $\ln v_d \propto \ln\dot\varepsilon_p$ and the slopes of the lines increased with rising temperature. We deduced from Eq. (10) that

$$\frac{d\ln v_d}{df}\approx\frac{A}{T}\left[\ln\frac{af^{1/2}}{T}+\frac{1}{2}\right]. \tag{13}$$

Since this derivative did not include any parameters of the centers (R_{1t}, n_t), it was interesting to compare the experimental values of this derivative and those deduced from Eq. (13). Table I gives the results of such a comparison. We can see from this table that the ag agreement between the experimental results and theory is good.

Figure 3 shows the temperature dependence of $f=qb/e$ We can see from Eq. (12) that the linear dependence $f=C_1-TC_2$ is obtained on the basis of our model if $\sum_t' n_t R_{1t}$ depends weakly (not exponentially) on T. This situation occurs if the E_t centers lie below the Fermi level μ and the position of the latter is independent of temperature (for example, in the case of strongly compensated semiconductors such as the samples employed in our study). The E_t centers lying above μ can, in this case, be ignored because of the low electron density n_t. However, if, because of the condition (11), a dislocation collects electrons from centers with $E_t>\mu(E_t-\mu\gg T)$, where $n_t\propto\exp[(\mu-E_t)/T]$, it then follows from Eq. (12) that $f(T)=\text{const}$, i.e., such centers can only give rise to a very weak (logarithmic) temperature dependence of f.

Finally, Fig. 4 shows the dependence f on the wavelength λ of light incident on a sample. It is clear from this figure that the illumination of ZnSe crystals increases f, as it does in the case of ZnS.[3] The maximum of the spectral dependence occurs at the fundamental absorption edge. Under these conditions, the maximum number of electron–hole pairs is created in the bulk of a sample. The resultant electrons are captured by the empty trapping levels [the value of $\sum_t' n_t R_{1t}$ in Eq. (12) increases]. It is less easy to explain why f is affected by light of long wavelengths λ right up to 8000 Å ($h\nu\approx 1.5$ eV) since the valence band–empty trapping level transitions should not occur beyond $h\nu\approx E_0-\mu\approx 2$ eV ($\lambda\approx 6200$ Å). The only optical transition which can increase f in this photon energy range seems to be the valence band–dislocation level transition. Therefore, we shall assume that the long-wavelength edge (1.5 eV) observed in the spectrum of Fig. 4 is associated with such transitions. Then,

$$E_d\approx-(E_G-1.5\,[\text{eV}])\approx-1.2\,[\text{eV}]. \tag{14}$$

Solving Eq. (10) for E_d and assuming the parameters $\omega_{dc}\sim E_d/\hbar\sim 10^{15}$ sec^{-1}, $v_d\sim 10^{-3}$ cm/sec, $\sum_t' n_t\sim 10^{16}$ cm^{-3} (deduced from the thermally stimulated conductivity and by the injected current method), $f=0.5$, and $T=0.025$ eV, we obtain $E_d\approx -1.3$ eV, which is very close to the value given by Eq. (14). In this estimate, the error in the value of $v_d\sum' n_t/\omega_{dc}$, amounting to two orders of magnitude can displace E_d by just 0.1 eV.

TABLE I.

T, K	$\left(\frac{d\ln v_d}{df}\right)_{\text{exp}}$	$\left(\frac{d\ln v_d}{df}\right)_{\text{theor}}$
299	87	78
306	81	75
385	42	45
434	29	33

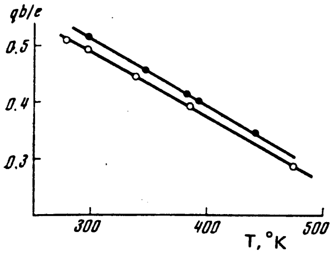

FIG. 3. Temperature dependences of the dislocation charge $qb/e=f$ in ZnSe for two samples deformed at a rate $\epsilon_p=10\mu$/min.

We shall now estimate the occupancy of a dislocation at rest with an energy level $E_d\approx -1.2$ eV, which is in thermal equilibrium with the lattice. It follows from Eq. (1) that

$$f=\left\{1+\exp\left[\frac{E_d+Af\ln(R_{\text{scr}}f/a)-\mu}{T}\right]\right\}. \tag{15}$$

If we assume that the screening radius R_{scr} is governed by the concentration of the E_t centers, i.e., $R_{\text{scr}}\sim(f/\pi n_t a)^{1/2}\sim 10^{-5}$ cm, we find that $f\sim 0.15$ at room temperature. If R_{scr} is taken to be the Debye radius or the distance between the dislocations, the value of f is found to be even smaller. Thus, a moving dislocation is characterized by a much higher electron occupancy than a dislocation at rest.

APPENDIX I

Equations (7) and (8) describe the variation of the electron density at impurities $n_t(x,y)$ with time or space. In view of the assumed homogeneous distribution of the impurity centers in space ($N_t=\text{const}$), the function n_t should be stationary in a coordinate system linked to the core of a moving dislocation (Fig. 1). The system (7)–(8) is easier to solve because transitions between the smooth parts of the distribution function occur at distances between an impurity center and a dislocation which are, on the one hand, much greater than the lattice constant and, on the other, much smaller than the screening radius $R_{\text{scr}}\sim 10^{-4}$ cm. Therefore, the electrostatic potential acting on an impurity because of the presence of a dislocation can be described by

$$\varphi=\begin{cases}\dfrac{2ef}{a\varepsilon}\ln\left(\dfrac{R_{\text{экр}}}{|x|}\right), & \dfrac{a}{f}\ll|x|\ll R_{\text{scr}}\\[2mm] \dfrac{2ef}{a\varepsilon}\ln\left(\dfrac{R_{\text{экр}}f}{a}\right), & |x|\leqslant\dfrac{a}{f}.\\[2mm] 0 & |x|>R_{\text{scr}}\end{cases} \tag{I.1}$$

The solution of Eq. (8) for an impurity center E_t is then

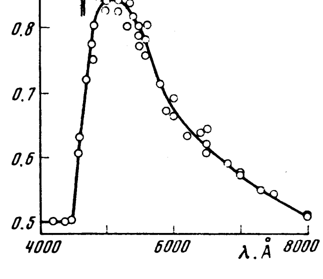

FIG. 4. Spectral dependence of the dislocation charge in ZnSe on the wavelength of incident light; $T=295$ °K, $\epsilon_p=10\,\mu$/min.

$$n_t(x,y)=n_t(-\infty)\exp\Big\{-\frac{\omega_{td}}{v_d}\int_{-\infty}^{x} g(x',y)$$

$$\Big[(1-f)+f\exp\Big(\frac{E_d-E_t(x',y)}{T}\Big)\Big]\,dx'\Big\}+\frac{N_t\omega_{td}}{v_d}\int_{-\infty}^{x} dx'\,g(x',y)f \tag{I.2}$$

$$\times\exp\Big\{-\frac{\omega_{td}}{v_d}\int_{x'}^{x} dx''g(x'',y)\Big[(1-f)+f\exp\Big(\frac{E_d-E_t(x'',y)}{T}\Big)\Big]\Big\}.$$

[We have considered here only the electron arrival and departure processes without a change in the coordinate z of an electron along a dislocation. This change can be considered but results in a slight renormalization of the order of $\ln^{1/2}(\omega_{td}a/\nu)$ of the transition attempt frequency ω_{td}.]

We shall simplify Eq. (I.2) as follows. We shall consider two characteristic distances of an impurity center from a dislocation, R_1 and R_2:

$$R_1=r_{0t}\ln\Big[\frac{R_1\omega_{td}(1-f)}{v_d}\Big], \tag{I.3}$$

$$R_2=\frac{a}{f}\exp\Big[\frac{E_t-E_d}{Af}\Big]. \tag{I.4}$$

Since the electron-filled centers become deep, $E_t \lesssim \mu \approx 0.7$ eV, the wave functions of the electrons at these centers are strongly localized and r_{0t} can be taken to be equal to the interatomic distance 2×10^{-8} cm. Then, for typical values $\omega_{td}\sim(E_1\sim E_d)/\hbar\sim10^{15}$ sec^{-1}, $v_d\sim10^{-3}$ cm/sec, and $f\approx0.5$, the distance R_1 is of the order of 5×10^{-7} cm, and it should be noted that an error in ω_{td} of five (!) orders of magnitude alters R_1 by a factor of less than 2. The distance R_2 becomes equal to R_1 for $E_t-E_d\approx0.6$ eV ($f\approx0.5$). The distance R_1 is such that, for this and shorter distances, the probability of electron loss from an impurity to a dislocation during the time that the dislocation travels one interatomic distance is very nearly equal to unity. This is justified if the impurity level obeys $E_t(R_1)\geq E_d$, i.e., if $R_1<R_2$. In the opposite case of $R_1>R_2$, the probability of electron loss from a dislocation to an empty level E_t becomes greater, beginning from distances R_1, and the level E_t is filled right up to distances R_2 because the loss of an electron from a dislocation to the corresponding center is preferred to the opposite process.

At distances shorter than R_2, the density of the electrons at the E_t center is low. We shall now consider the situation at the tail end of a dislocation. If $R_1<R_2$, the flux from a dislocation to an impurity exceeds the reverse flux, beginning only from distances greater than R_2. However, at these distances, the probability of electron transfer from a dislocation to an impurity is exponentially small, $\sim\exp(-R/r_{0t})$. Consequently, the impurity centers move away from a dislocation in the empty state. The reverse is true for $R_1>R_2$. Then, beginning from distances R_2, the impurity centers become filled with electrons from a dislocation and they move away from it in the filled state. A graphical solution of Eq. (I.2) is shown in Figs. 5a and 5b. It is worth noting another interesting detail of this solution: transition from one smooth part of the distribution to another occurs over distances of the order of the interatomic separation. In the case of $R_1<R_2$, this is self-evident be-

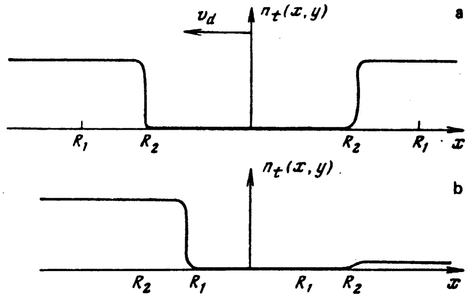

FIG. 5. Schematic diagram illustrating the solution of Eq. (8): (a) case $R_1>R_2$; (b) case $R_1<R_2$.

cause the frequency of transitions per unit time is proportional to $\omega_{td}\exp(-R/r_{0t})$ and a change in the distance by a lattice constant alters this frequency approximately by a factor of e. In the opposite case of $R_2<R_1$, the argument of the exponential function changes by a factor e when the distance is altered by R_2T/Af, which is again of the order of the interatomic separation.

We have considered so far only the centers with a definite energy E_t. In a real semiconductor, the distribution of the energies of impurity centers may be almost continuous. We then obtain an interesting situation. For small differences of the impurity level energies E_t from the dislocation value E_d, the distance R_2 is small and an impurity which temporarily gives up an electron to a dislocation recovers this electron after the dislocation passes close to it. In the case of centers with higher energies, the inequality $R_1>R_2$ is not obeyed because the distance R_1 is practically independent of the level position and R_2 is an exponential function of this position, in accordance with Eq. (I.4). The difference between the energies beginning from which an impurity center gives up its electron to a dislocation can be found from

$$E_t-E_d=Af\ln\Big\{\frac{r_{0t}}{a}f\ln\Big(\frac{\omega_{td}r_{0t}}{v_d}\Big)\Big\}\sim Af\approx0.4\text{ eV}\,(f\approx0.5). \tag{I.5}$$

Finally, the flux of electrons from the E_t centers to a dislocation (per unit time and unit dislocation length) is

$$(1-f)v_d\Big({\sum_t}' n_tR_{1t}\Big), \tag{I.6}$$

where the summation is carried out over all the centers E_t satisfying

$$E_t-E_d\geqslant Af\ln\Big(\frac{r_{0t}f}{a}\ln\Big(\frac{\omega_{td}r_{0t}}{v_d}\Big)\Big). \tag{I.7}$$

Equating the flux of electrons from the E_t centers to a dislocation to the flux of electrons from the dislocation to the conduction band [Eq. (4)], we obtain the following equation for the occupancy f:

$$(1-f)v_d\Big({\sum_t}' n_tR_{1t}\Big)=\omega_{dc}\Big(\frac{\pi e^2f^3}{T\varepsilon a^3}\Big)^{1/2}\exp\Big\{\frac{E_d+Af[\ln(a'f^{1/2}/T)-1]}{T}\Big\}. \tag{I.8}$$

APPENDIX II

We have so far ignored the probability of electron loss from a local level E_t to the conduction band when the center with this level is approached by a dislocation to the shorter of the two distances R_1 or R_2. This has made it possible to assume that, before the interaction with a dislocation, these centers have the electron density $n_t(-\infty)$. We shall justify this assumption by estimating the probability of an electron transition from a level E_t to the conduction band when the distance between the center with this level and a dislocation varies from R_{scr} to R_1. In such estimates, the important range of distances is that exceeding R_1 but smaller than a certain value R_3. Beginning from R_3, an electron from the level E_t is transferred to the conduction band not at a given point in space but as a result of tunneling over a certain distance. An analysis similar to that used in calculations of the tunneling of an electron from a dislocation to the conduction band[14] gives the following result:

$$P(x, E_c-E_t)=P_0(x, E_c-E_t)\exp\left\{\frac{Af}{T}\left[\ln\frac{R_3}{x}-1+\frac{x}{R_3}\right]\right\}, \quad \text{(II.1)}$$

where $P(x, E_c-E_t)$ is the probability of a transition (per unit time) of an electron from E_t to the conduction band in the field of a dislocation at a point x, P_0 is the same probability but outside the dislocation field, $R_3=r_0Af/T$, where r_0 is the order of the interatomic distance. The quantity in the argument of the exponential function in Eq. (II.1) is of the order of $(R_3-x)^2/R_3^2$ when the difference between x and R_3 is small. The values of R_1 and R_3 are of the same order of magnitude and, even if they differ by a factor of 3, we have

$$\exp\left\{\frac{Af}{T}\left[\ln\frac{R_3}{R_1}-1+\frac{R_1}{R_3}\right)\right\}\approx\exp\left\{\frac{0.1[\mathrm{eV}]}{T'}\right\}. \quad \text{(II.2)}$$

The probability of thermal release of an electron from a level in a time dt is:[15]

$$P_0\,dt=\frac{\langle v\sigma\rangle N_c}{g}\exp\left\{\frac{E_t-E_c}{T}\right\}dt, \quad \text{(II.3)}$$

where v is the thermal velocity of the electrons, σ is the electron-capture cross section of an E_t center, N_c is the effective density of states in the conduction band, and $g=2$. Without allowance for the tunneling in the dislocation field (at room temperature), we have

$$P_0 dt\sim\frac{10^7\cdot10^{-15}\cdot2\cdot10^{18}}{2}\exp\left\{-\frac{0.7}{0.025}\right\}\frac{10^{-4}}{10^{-3}}\sim10^{-3}\ll1,$$

$$dt\sim\frac{R_{scr}}{v_d}\sim\frac{(N_d)^{-1/3}}{v_d}\sim\frac{10^{-4}}{10^{-3}}\sim10^{-1}\,\mathrm{sec}.$$

The possibility of an indirect transition in the electric field of a dislocation at distances R_3-R_1 contributes the following amount to this probability:

$$dP'\sim P_0(E_t-E_c-0.1\ \mathrm{eV})\frac{R_3-R_1}{v_d}\sim10^{-3}\ll1,$$

i.e., the probability of thermal release of an electron from filled levels ($E_t\lesssim\mu\sim0.7$ eV) is negligible up to the moment when a given center is emptied because of the transfer of an electron to a dislocation. Therefore, we may assume that, up to distances of the order of 10^{-6} cm, we have $n_t(x,y)=n_t(-\infty)$.

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