Cooling of an electron gas in many-valley semiconductors in the case of inelastic carrier scattering

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An investigation is reported of the cooling of an electron gas (reduction in the average energy of electrons on increase in the electric field E) in many-valley semiconductors which occurs when the dominant energy dissipation mechanism is the generation of phonons with the Debye temperature $\hbar\omega/k_0$ much higher than the lattice temperature T and the dissipation of momentum is primarily due to the scattering by ionized impurities. It is shown that the impurity intervalley scattering time $\tau_{iv}^{(I)}$, which increases on increase in E in the case of carrier heating, decreases on increase in E in the cooling case irrespective of whether intravalley $\hbar\omega_0$ or intervalley $\hbar\omega_{iv}$ phonons are generated. The phonon intervalley scattering time $\tau_{iv}^{(ph)}$ decreases on increase in E in the case of heating and continues to decrease when carriers are cooled if the cooling is due to generation of intevalley phonons $\hbar\omega_{iv}$ that determine $\tau_{iv}^{(ph)}$. However, if the cooling is due to generation of intravalley phonons of energy $\hbar\omega_0 < \hbar\omega_{iv}$, then $\tau_{iv}^{(ph)}$ rises on increase in the field. A change in the sign of the derivative $d\tau_{iv}/dE$ on transition from carrier heating to cooling causes an inversion of the valley population. It is shown that an investigation of the intervalley population transfer and of the electrical conductivity anisotropy can provide reliable diagnostic tools for the study of the cooling of carriers in semiconductors.

1. INTRODUCTION

Cooling of carriers in semiconductors was predicted by Gribnikov and Kochelap[1] and observed experimentally by Ashmontas et al.[2,3] It involves a reduction in the total energy of the electron gas in a semiconductor on increase in the electric field when the dominant energy dissipation mechanism is the generation of optical phonons and the momentum is dissipated primarily by scattering on ionized impurities. An electron diffusing in an electric field attains the energy of an optical phonon $\hbar\omega_0$, (the constant for the interaction with such a phonon is assumed to be large) and after emitting this phonon drops to the range of low energies from where escape by acquisition of energy from the electric field is difficult because of the high rate of impurity scattering, the probability of which is $\sim\varepsilon^{-3/2}$ (ε is the carrier energy). If the electron flux to the range of low energies because of such optical scattering is sufficiently large, the majority of electrons become localized in this energy range and this lowers their average energy $\bar{\varepsilon}$ ($\bar{\varepsilon}$ can fall to a value smaller than the equilibrium value $3k_0T/2$, which is in particular observed in Refs. 2 and 3). The cooling effect occurs in the temperature range $k_0T \ll \hbar\omega_0$ at extremely low carrier densities, when their interaction with one another can be ignored. The range of electric fields in which this effect occurs was considered in Ref. 1.

Further studies of the effects associated with the cooling of carriers was made by one of the present authors[4] and it was found that ionization of carriers from impurities gives rise to a negative differential conductance. The magnetoresistance and field dependences of the carrier mobility under cooling conditions in III-V semiconductors were calculated recently[5] and an analysis was made of the available experimental data on these dependences.

We shall consider theoretically the effects associated with the cooling of electrons in many-valley semiconductors. We shall consider two types of semiconductor. Type A is a semiconductor such as n-type Ge in which the energy of a typical intervalley phonon $\hbar\omega_{iv}$ is small compared with the energy of an intravalley optical phonon $\hbar\omega_0$ (Ref. 6) and, moreover, the constant W_{iv} of the interaction with an intervalley phonon is small compared with the constant W_0 of the interaction with an optical phonon, so that intervalley scattering has practically no influence on the energy dependence of the distribution function. Type B represents semiconductors such as silicon or n-type diamond, when the interaction with an intervalley phonon is characterized by a large constant W_{iv} (Ref. 7) and it is this phonon that dominates the energy scattering in the range of electric fields of interest to us.

2. INELASTIC SCATTERING OF ELECTRONS BY INTRAVALLEY OPTICAL PHONONS

In type A semiconductors the weak interaction with intervalley phonons means that the condition of independent energy balance of the valleys is obeyed and in calculating the distribution functions of the valleys we can ignore the intervalley scattering process.[8]

In Ge, Si, and other semiconductors the anisotropy of the relaxation times of the electron momentum dissipated on intravalley acoustic phonons is weak compared with the anisotropy of the effective masses.[8,9] Therefore, if we assume that the scattering is isotropic, we find that the application of the Herring–Vogt transformation,[10] which converts ellipsoidal constant-energy surfaces to the spherical shape, makes the transport equations for electrons in different valleys identical apart from the value of the effective electric field in the valleys:

$$E_i^* = E\left[\frac{m}{m_\parallel}\cos^2\varphi_i + \frac{m}{m_\perp}\sin^2\varphi_i\right]^{1/2}. \tag{1}$$

Here, ϕ_i is the angle between the longitudinal axis of the constant-energy ellipsoid for the i-th valley and the direction of the total electric field E; $m = (m_\parallel m_\perp^2)^{1/3}$ is the density-of-states effective mass; $m_\parallel^{-1}$ and $m_\perp^{-1}$ are the longitudinal and transverse

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components of the tensor of the reciprocal effective masses. The equations for the transformation of the variables, which we shall now identify by an asterisk, are the same as the equations for a semiconductor with an isotropic dispersion law.[1,7,11]

We shall adopt the quasielastic approximation in the passive range of energies $\varepsilon < \hbar\omega_0$ (Ref. 12) and describe the distribution function $F_{\mathbf{p}^*}^{(i)}$ of of the electron momentum $\mathbf{p}^*$ in the i-th valley by the expression

$$F_{\mathbf{p}^*}^{(i)} = F_0^{(i)}(\varepsilon) + F_1^{(i)}(\mathbf{p}^*), \quad F_1^{(i)}(\mathbf{p}^*) \ll F_0^{(i)}(\varepsilon). \tag{2}$$

The isotropic $F_0^{(i)}$ and anisotropic $F_1^{(i)}(\mathbf{p}^*)$ parts of this distribution function are found from

$$x^2[\Theta_i(x)+1]\frac{dF_0^{(i)}(x)}{dx} + x^2F_0^{(i)}(x) = -J_i(x), \tag{3}$$

$$F_1^{(i)}(\mathbf{p}^*) = \frac{e\tau_{\mathbf{p}^*}(x)}{k_0T}(\mathbf{E}_i^*\mathbf{V}^*)\frac{dF_0^{(i)}(x)}{dx}, \tag{4}$$

where

$$\Theta_i(x) = \mathcal{E}_i^2\frac{x}{x^2+B}, \quad x = \frac{\varepsilon}{k_0T}, \quad B = 6\frac{\mu_a^*}{\mu_I^*}. \tag{5}$$

$$\mathcal{E}_i = \frac{E_i^*}{E_0}, \quad E_0 = \frac{4s}{\mu_a^*}\sqrt{\frac{1}{3\pi}\frac{m_\perp}{m}\frac{\Sigma_0^2}{\Sigma_1^2}}. \tag{6}$$

Here, s is the longitudinal velocity of sound; Σ_0^2 and Σ_1^2 are the deformation potential constants[9-11]; e is the electron charge; $\tau_{\mathbf{p}^*}(x)$ is the momentum relaxation time. It is assumed that the electron momentum dissipated on acoustic intravalley phonons and ionized impurities, and that $\tau_{\mathbf{p}^*}$ is described by

$$\tau_{\mathbf{p}^*}(x) = \tau_0\frac{x^{3/2}}{x^2+B}, \quad \tau_0 = \frac{3\sqrt{\pi}}{4e}m\mu_a^*. \tag{7}$$

The constant B in Eqs. (5) and (7) allows for the contribution of each of these mechanisms and can be expressed in terms of the low-field mobilities μ_a^* and μ_I^* of Eq. (5), which represent the scattering solely by acoustic phonons or solely by ionized impurities. It should be pointed out that the anisotropy of the momentum relaxation time due to the scattering by impurities in n-type Ge and n-type Si is of the order of the anisotropy of the effective masses[13] so that the effective field method cannot be used when several scattering mechanisms act simultaneously.[8] However, when one of these mechanisms predominates, which will be the case of interest to us, an allowance for the anisotropy of $\tau_{\mathbf{p}^*}$ alters only the coefficients in Eq. (1).

The right-hand side of Eq. (3) allows for relaxation of the distribution function because of the interaction of electrons with optical phonons, whereas $J_i(x)$ is equal, apart from a constant, to the flux of electrons in the one-dimensional energy space associated with this interaction:

$$J_i(x) = \frac{g(k_0T)k_0T}{\tau_{\omega_0}(k_0T)}\int_0^x u^{1/2}\sqrt{u+x_0}\,[(N_0+1)F_0^{(i)}(u+x_0) - N_0F_0^{(i)}(u)]\,du, \tag{8}$$

$$J_i(x) = \frac{\tau_\varepsilon(k_0T)}{g(k_0T)k_0T}j_i(x), \quad 0 \leqslant x \leqslant x_0, \quad x_0 = \frac{\hbar\omega_0}{k_0T}, \tag{9}$$

where N_0 is the equilibrium distribution function of optical phonons; $g(k_0T) = (2m^3k_0T)^{1/2}/\pi^2\hbar^3$ is the density of states at $\varepsilon = k_0T$; $\tau_\varepsilon(k_0T) = \tau_0(\Sigma_1^2/\Sigma_0^2)\cdot(k_0T/2m_\perp s^2)$ and $\tau_{\omega_0}(k_0T) = \sqrt{2}\,\pi\hbar^3\rho\omega_0/W_0^2m^{3/2}(k_0T)^{1/2}$ are the characteristic relaxation times of the electron energy and of the optical phonon emission, respectively; ρ is the density of the investigated crystal.

In a fairly wide range of fields the penetration of electrons to the active range of energies $\varepsilon > \hbar\omega_0$ is limited to a narrow interval $\Delta\varepsilon \ll \hbar\omega_0$, so that we can quite accurately regard the fluxes of Eq. (8) and the quantities of Eq. (9) as constants[1,8] and assume that the boundary condition for Eq. (3) is the approximate expression $F_0^{(i)}(x_0) \simeq 0$ (Ref. 1).

Under these conditions we can use directly the solution of the transport equation (3) obtained in Ref. 1 for a single-valley isotropic semiconductor

$$F_0^{(i)}(x) = J_i\exp[-I_i(x)]\int_x^{x_0}\frac{\exp[I_i(u)-I_i(x_0)]}{u^2[\Theta_i(u)+1]}du, \tag{10}$$

$$I_i(x) = \int_0^x\frac{du}{\Theta_i(u)+1}.$$

In Eq. (10) the expression $\bar{J}_i = J_i\exp[I_i(x_0)]$ represents the constant of integration, calculated from the conditions of normalization of the distribution functions $F_0^{(i)}(x)$ to the electron density n_i in each of the valleys, which is given by

$$n_i = \frac{2}{(2\pi\hbar)^3}\int F_0^{(i)}(x)\,d\mathbf{p}^*. \tag{11}$$

The density n_i is in its turn found from the conditions of electrical neutrality and balance of intervalley transitions, which give[8]

$$n_i = n\frac{\tau_i}{\sum_{k=1}^{\lambda}\tau_k}, \tag{12}$$

where τ_i is the time for the intervalley scattering of an electron from a valley i to any other valley, λ is the number of valleys, and n is the total density of electrons in a semiconductor.

The distribution function $F_{\mathbf{p}^*}^{(i)}$ in a valley i depends only on the effective field $\mathcal{E}_i$, so that the field dependences of the average energy $\bar{\varepsilon}_i$ and of the drift velocities $\bar{v}_i$ in the valleys are the same as in the case considered in Ref. 1. Therefore, we shall consider only those effects which are associated with the field dependences of τ_i. The intervalley scattering time τ_i can be calculated from Eq. (10):

$$\frac{1}{\tau_i} = \int_0^{x_0}\frac{x^{1/2}}{\tau_{iv}(x)}F_0^{(i)}(x)\,dx \bigg/ \int_0^{x_0}x^{1/2}F_0^{(i)}(x)\,dx, \tag{13}$$

where the frequencies of intervalley transitions $\tau_{iv}^{-1}(x)$ should allow for the contribution of the scattering by phonons and by ionized impurities:

$$\frac{1}{\tau_{iv}(x)} = \frac{1}{\tau_{iv}^{(ph)}(x)} + \frac{1}{\tau_{iv}^{(I)}(x)} = (\lambda-1)\left[\frac{\sqrt{x-x_{iv}}}{\tau_{iv_0}^{(ph)}} + \frac{x^{-r}}{\tau_{iv_0}^{(I)}}\right], \quad x_{iv} = \frac{\hbar\omega_{iv}}{k_0T}. \tag{14}$$

In the intervalley transfer time due to scattering by phonons $\tau_{iv}^{(ph)}(x)$ only the processes of the scattering accompanied by the emission of a phonon of energy $\hbar\omega_{iv}$ are included; $\tau_{iv_0}^{(ph)} = \frac{\sqrt{2}\,\pi\hbar^3\rho\omega_{iv}}{W_{iv}^2m^{3/2}(k_0T)^{1/2}}$, the intervalley impurity scattering is assumed to be characterized by a power-law dependence of $\tau_{iv}^{(I)}$ on x (Refs. 14 and 15) ($\tau_{iv_0}^{(I)}$ is a constant dependent on the impurity concentration and r is a parameter).

It is clear from Eq. (12) that the intervalley population transfer is governed entirely by the times τ_i and the highest population is exhibited by that valley for which τ_i is longest. In fields such that the condition $I_i(x_0) \gg 1$ is satisfied, this means physically that the flux of Eq. (8) is small, that the main contribution to the integral (10) comes from the upper limit, and that integral is practically constant. The distribution function then has the form $F_0^{(i)}(x) \approx C_i \exp[-I_i(x)]$, where C_i is the integration constant, relabeled compared with Ref. 10. The effects associated with intervalley redistribution in this range of fields were studied in detail in Refs. 8 and 16. Here, τ_i rises exponentially on increase in $\mathcal{E}_i$ when the intervalley phonon scattering predominates and this results in emptying of the valleys with higher values of $\mathcal{E}_i$, so that inverse Sasaki effect is exhibited by the valley populations.[8] If intervalley impurity scattering takes place, it may predominate in weak fields, giving rise to the anomalous Sasaki effect[16] when there are more electrons in a valley with a higher value of $\mathcal{E}_i$. When the field $\mathcal{E}_i$ is increased, the anomalous Sasaki effect changes to the normal effect and this occurs in fields which are still such that the condition $I_i(x_0) \gg 1$ is obeyed.

In the range of fields where the condition $I_i(x_0) \gg 1$ is satisfied, the main dependence of the distribution functions on the energy in Eq. (10) is determined by the lower limit of the integral and these distributions are described by

$$F_0^{(i)}(x) = J_i \int_x^{x_0} \frac{du}{u^2[\Theta_i(u)+1]}. \qquad (15)$$

We shall consider separately two cases of predominance of specific mechanisms in the dissipation of electron momentum.

a) The electron momentum relaxes by interaction with acoustic phonons ($B \ll \bar{x}$). It then follows from the condition $I_i(x_0) \ll 1$ that $\Theta_i(x) = \mathcal{E}_i^2/x \gg 1$ and the distribution function in the range of energies of interest to us is independent of the field.[1] According to Eq. (13) the field dependence τ_i for a valley satisfying the condition $I_i(x_0) \ll 1$ then reaches saturation. This means that an increase in the electric field weakens the intervalley distribution which occurs in the range of fields where $I_i(x_0) \gg 1$ and when all the fields satisfying the condition $I_i(x_0) \ll 1$ ($i = 1, 2, \ldots, \lambda$), the redistribution ceases. Consequently, in this range of fields the Sasaki effect affecting the redistribution between the valleys is no longer observed. The average energy then also reaches saturation and becomes $\varepsilon = {}^9/_{25}\,\hbar\omega_0$.

b) We shall now consider the case when the electron momentum is dissipated by ionized impurities ($B \gg x_0$). In this case we have $\Theta_i(x) = (\mathcal{E}_i^2/B)$ and, in contrast to the preceding case, the condition $I_i(x_0) \ll 1$ does not imply that $\Theta_i(\bar{x}) \gg 1$ is obeyed. This is due to a reduction in $\Theta_i(\bar{x})$ on increase in $\bar{x}$ and is the reason for the cooling of carriers.[1] Calculation of the average energy with the aid of Eq. (15) gives $\varepsilon_i = (6\sqrt{Bx_0}/5\pi\mathcal{E}_i)\,k_0T$, i.e., this energy decreases on increase in the field and such cooling of electrons begins earlier in a valley with the highest effective field $\mathcal{E}_i$, i.e., in the valley which is heated most strongly cooling begins earlier than else where. In the case of τ_i it follows from Eqs. (13)-(15) that

$$\frac{1}{\tau_i} = \frac{\alpha}{\tau_{iv_0}^{(ph)}}\,\mathcal{E}_i^{-1} + \frac{\beta}{\tau_{iv_0}^{(I)}}\,\mathcal{E}_i^{2r}, \qquad (16)$$

where

$$\alpha = \frac{3\sqrt{B}}{4\pi}\int_\delta^1 t^{1/2}(t-\delta)^{1/2}(t^{-2}-1)\,dt, \quad \beta = \frac{3}{(3-2r)\,B^r\cos\pi r}, \quad \delta = \frac{x_{iv}}{x_0}.$$

In the calculation of Eq. (16) it is assumed that $r < {}^1/_2$. However, if $r \geq {}^1/_2$, when the corresponding integrals diverge logarithmically or in accordance with a power law so that to remove this divergence it is necessary to allow for finite penetration of electrons to the active range of energies or to use the more correct dependence $\tau_M^{(I)}(x)$ in the limit $x \to 0$. It follows from Eq. (16) that if

$$\frac{\alpha}{\beta}\frac{\tau_{iv_0}^{(I)}}{\tau_{iv_0}^{(ph)}} \gg \mathcal{E}_i^{2r+1}, \qquad (17)$$

then $\tau_i \sim \mathcal{E}_i$, i.e., when the phonon mechanism predominates in the intervalley scattering the population of a valley with a higher field $\mathcal{E}_i$ (which is the valley that becomes cooler as a result of the cooling process) is higher. Consequently, if the condition $I_i(x_0) \ll 1$ is satisfied by all the valleys, then the anomalous Sasaki effect applies to the valley populations. It should be stressed that in contrast to Ref. 16 this occurs in a different range of fields and for the phonon mechanism of intervalley scattering (which is an important qualification). In addition to Eq. (17), we can expect the condition $B \gg \bar{x}$ to be obeyed, whereas the strong inequality $B \gg x_0$ need not be obeyed. If the impurity concentration is such that in the range of energies $\varepsilon \sim \hbar\omega_{iv}$ the acoustic momentum scattering predominates, we still have τ_i described by Eq. (16), but with a somewhat different value of α. The possibility that Eq. (17) is satisfied is supported also by numerical results from Ref. 15.

If the condition opposite to Eq. (17) is obeyed, then $\tau_i \sim \mathcal{E}_i^{-2r}$ and therefore in the case of predominance of the impurity intervalley scattering the valley with a high effective field is less populated (normal Sasaki effect), whereas in an analogous situation in the case of weak heating we have the anomalous Sasaki effect.[16] The physical reason for this is the fact that a valley with a higher effective field $\mathcal{E}_i$ is colder, i.e., a form of inversion of the valleys in respect of the degree of heating and population takes place.

We have discussed the Sasaki effect only in the case of intervalley redistribution of electrons. If a sample is bounded in the transverse direction, a Sasaki field $E_\perp$ appears in this direction and it can be calculated from the absence of the total current along this field. For the sake of simplicity, we shall consider the appearance of the Sasaki field in a two-valley semiconductor when the valleys are oriented at 90° relative to one another and their symmetry axis makes an angle of ψ with the direction of the applied external field E_Σ. Using Eqs. (1), (4), and (12), we find by analogy with Ref. 8 that the Sasaki field $E_\perp$ is described by

$$E_\perp = aE_\Sigma\frac{F\cos 2\psi}{1+aF\sin 2\psi}, \quad F = \frac{K-1}{K+1}, \qquad (18)$$

where $a = (m_\parallel - m_\perp)/(m_\parallel + m_\perp)$ is the anisotropy parameter and K is the ratio of the products of the mobility in the valleys and the time of escape from the valleys, which can be represented in the form

$$K = \frac{\int\limits_0^{x_0} \frac{x^3}{x^2+B} \frac{dF_0^{(2)}(x)}{dx}\,dx \Big/ \int\limits_0^{x_0} \frac{x^{1/2}}{\tau_{iv}(x)} F_0^{(2)}(x)\,dx}{\int\limits_0^{x_0} \frac{x^3}{x^2+B} \frac{dF_0^{(1)}(x)}{dx} \Big/ \int\limits_0^{x_0} \frac{x^{1/2}}{\tau_{iv}(x)} F_0^{(1)}(x)\,dx} \tag{19}$$

The sign of the Sasaki field for a given orientation of the valleys is governed by the sign of F. We shall consider the specific case $\mathscr{E}_1 > \mathscr{E}_2$, and, therefore, we shall select $\psi > 0$ and label 1 the valley which is characterized by $\phi_1 = 45° + \psi$ and label 2 the valley which satisfies $\phi_2 = 45° - \psi$ [see Eq. (1)].

In the case (a) if $I_i(x_0) \ll 1$ (i = 1, 2) the drift velocities in the valleys deduced from Eqs. (4) and (15) are proportional to the external field E_X ("second" ohmic region[1,5]), and it follows from Eqs. (18) and (19) that $E_\perp = 0$. Consequently, in this range of fields for any orientation of the valleys there is no Sasaki effect in the intervalley redistribution or in the appearance of the transverse field, and a semiconductor can be regarded as isotropic. An increase in the electric field causes electrons to penetrate deeply to the active range of energies so that the above analysis is no longer valid, but it is clear from qualitative considerations that the ordinary Sasaki effect appears in such fields. Therefore, the dependences on the applied field E_X have two extrema both in the case of the redistribution of valleys and in the case of the transverse field $E_\perp$. The first extremum occurs in the range of fields corresponding to the change from the condition $I_i(x_0) \gg 1$ to the opposite one for the valley with the lowest effective field $\mathscr{E}_i$ whereas the second extremum occurs outside the range of strong electric fields of interest to us and it is discussed in Ref. 8.

In the case (b) the drift velocities of the valleys reach saturation and the behavior of the Sasaki field $E_\perp$ depends on whether the condition (17) is obeyed. If it is obeyed, it then follows from Eq. (19) that K = 1 and $E_\perp = 0$, i.e., in spite of the anomalous intervalley redistribution, the Sasaki field is absent for any orientation of the valleys and this is true of a semiconductor with an arbitrary number of valleys provided the valley with the lowest effective field $\mathscr{E}_i$, satisfies not only Eq. (17), but also the condition $I_i(x_0) \ll 1$.

When the condition opposite to Eq. (17) is obeyed, it follows from Eqs. (18) and (19) that $K = (\mathscr{E}_1/\mathscr{E}_2)^{2r+1} \gg 1$ and $F \approx 1$, i.e., the Sasaki field differs from zero and its sign corresponds to the normal effect (and this is true also of the populations of the valleys). It should be stressed that the Sasaki effect is realized for intervalley impurity scattering, when the anomalous Sasaki effect applies in weak fields,[16] i.e., an increase in the field E_X on transition from $I_i(x_0) \gg 1$ to the opposite case reverses the sign of $E_\perp$. It should also be pointed out that when the intervalley scattering is controlled by impurities the relationship between $\hbar\omega_0$ and $\hbar\omega$ is now of no importance.

3. INELASTIC SCATTERING OF ELECTRONS BY INTERVALLEY PHONONS

In type B semiconductors after inelastic scattering of a phonon of energy $\hbar\omega$ an electron is transferred to a different valley and we now have the case of an energy-dependent balance of the valleys,[8] when such scattering must be included in the transport Eq. (3). If in the balance between the intervalley fluxes we allow only for the phonon scattering, then the problem is of the kind discussed in Ref. 8, where is it shown that $j_i = j_0$ = const for all the valleys and if $I_i(x_0) \ll 1$, then the intervalley scattering time continues to decrease on increase in the field, but now in accordance with a power law ($\tau_i \sim \mathscr{E}_i^{-2}$) and throughout the range of fields in question the normal Sasaki effect applies.

If the electron momentum is dissipated on impurities, then the balance of the intervalley fluxes should allow not only for the phonon mechanism, but also for the intervalley impurity scattering. This results in renormalization of the fluxes j_i of Eqs. (3), (8), and (9), and we have $j_i \neq j_0$. On the right-hand side of Eq. (3) we find that in the case of the i-th valley we have the sum of arrivals from the active regions of other valleys and instead of J_i we must now write $\sum\limits_{k \neq i} J_{ki}$ [this substitution could also be made in Eqs. (10) and (15)] and the time of escape from the valley i is

$$\frac{1}{\tau_i} = \left[\frac{1}{\tau_s(k_0T)} \sum_{\substack{k \neq i \\ k=1}}^{\lambda} J_{ik} + \frac{1}{\tau_{x_0}^{(I)}} \int\limits_0^{x_M} x^{1/2-r} F_0^{(i)}(x)\,dx \right] \Big/ \int\limits_0^{x_M} x^{1/2} F_0^{(i)}(x)\,dx, \tag{20}$$

where the constant $\bar{J}_{ik}$ governing the flux from the active region of a valley i to the passive region of a valley k are found form the conditions of normalization and the intervalley balance. In particular, in the two-valley case ($\lambda = 2$) we find from Eqs. (15) and (20) that

$$\frac{1}{\tau_i} = \left[\frac{3}{2\pi} \frac{\nu_i/\nu_k}{\tau_s(k_0T)} \mathscr{E}_i + \frac{\beta}{\tau_{iv_0}^{(I)}} \mathscr{E}_i^{2r} \right], \quad i, k = 1, 2, \; i \neq k, \tag{21}$$

where

$$\nu_i = \frac{1}{\tau_s(k_0T)} - \frac{2\pi}{3} \beta \frac{1}{\tau_{iv_0}^{(I)}} \mathscr{E}_i^{2r-1}.$$

The second term in Eq. (21) is small compared with the first since the intervalley impurity scattering is ignored in the transport equation and its field dependence is the same as for type A semiconductors [Eq. (16)]. A special feature of Eq. (21) is that allowance for the impurity intervalley scattering renormalizes the intervalley phonon scattering time ν_i/ν_k). Therefore, in this case τ_i also decreases on increase in $\mathscr{E}_i$ and the allowance for the impurity intervalley scattering reduces, according to Eq. (21), the value of τ_i. Since on increase in $\mathscr{E}_i$ the process of cooling increases the number of electrons at low energies, this results in a stronger dependence of τ_i on $\mathscr{E}_i$.

We conclude by noting that in the cooling case the time for the intervalley scattering by impurities decreases on increase in $\mathscr{E}_i$ ($\tau_i \sim \mathscr{E}_i^{-2r}$), and then the mobility obeys $\mu_i \sim \mathscr{E}_i^{-1}$, i.e., it also decreases on increase in $\mathscr{E}_i$. Therefore, in the region of predominant impurity scattering in some semiconductors we can expect the criterion

(23.7) of Eq. (8) to be satisfied and multivalued distributions of electrons between the valleys may be realized, although it has been assumed earlier[8] that they occur only in the case of the phonon mechanism of intervalley scattering that ensures a reduction in τ_i on increase in the field.

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