Momentum relaxation time and temperature dependence of electron mobility in semiconductor superlattices consisting of weakly interacting quantum wells

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(Submitted January 22, 1999; accepted for publication March 23, 1999)
Fiz. Tekh. Poluprovodn. **33**, 1240–1245 (October 1999)

Formulas are derived for the longitudinal and transverse electron momentum relaxation times in a superlattice consisting of weakly interacting quantum wells in the approximation of a bulk-semiconductor scattering potential. Scattering by impurity ions, neutral atoms, and volume-type longitudinal acoustic and polar optical phonons is studied. A numerical analysis is performed of the longitudinal and transverse momentum relaxation times for scattering by impurity ions and neutral atoms as a function of the electron energy, temperature, density of the electron gas, and quantum well width. © *1999 American Institute of Physics.* [S1063-7826(99)02010-4]

As is well known, scattering mechanisms play an important role in the electric and optical properties of semiconductors and structures based on them. The present paper is devoted to the calculation and numerical analysis of the tensor of the reciprocal momentum relaxation time and mobility of electrons in semiconductor superlattices (SLs) consisting of weakly interacting quantum wells (QWs) based on heterostructures. Such SLs are of interest because of the unique properties of the electron energy spectrum and the possibility of using the SLs in infrared-range photodetectors[1–5] and microwave generators.[6–9] Formulas are derived for calculating the longitudinal and transverse electron momentum relaxation times in such SLs for the basic scattering mechanisms, and a numerical analysis is performed of the energy dependence of the relaxation time and temperature dependence of the mobility for scattering by impurity ions and neutral atoms. The following basic approximations were made in deriving the formulas: The electron gas in the quantum wells is almost two-dimensional, i.e., the width of the lower conduction miniband is much less than the average electron energy, the wave function of the lower miniband can be represented as a Bloch sum over wave functions of the ground state of infinitely deep QWs, and the potentials of the scattering centers in bulk semiconductors and SLs are essentially the same. These approximations together with the results obtaind by other authors[10–14] show that the formulas derived are acceptable not only for qualitative but also quantitative analysis. The latter is valid if the period of the SL is sufficiently large and size quantization of the phonon spectrum can be disregarded[12,14] and if the anisotropy of the scattering potentials, which arises primarily as a result of the difference in the dielectric constants of the layers of the SL which correspond to QWs and barriers, can also be disregarded.[11]

1. ELECTRON MOMENTUM RELAXATION TIME TENSOR

The tensor of the reciprocal of the relaxation time for electrons in the bottom miniband was calculated using the nonequilibrium correction $f_{1\mathbf{k}}$ for the equilibrium distribution function f_0 given by

$$f_{1\mathbf{k}}=e\frac{\partial f_0}{\partial E}\sum_i \tau_i F_i v_i, \tag{1}$$

where τ_i are the diagonal components of the relaxation time tensor, F_i and v_i are, respectively, the electric field intensity and electron velocity vectors in a Cartesian coordinate system, and E is the electron energy. As is well-known, for composite SLs consisting of weakly interacting QWs

$$E=\frac{\hbar^2(k_x^2+k_y^2)}{2m^*}+\frac{\Delta}{2}(1-\cos qd), \tag{2}$$

where m^* is the effective mass corresponding to the free motion of electrons along the layers of the SL, Δ is the width of the bottom miniband, corresponding to electron motion along the axis of the SL parallel to the z axis, $d=a+b$ is the SL period, a and b are the QW and barrier widths, respectively, and k_x, k_y, and $q=k_z$ are the components of the wave vector. In the two-dimensional electron gas approximation ($\Delta \ll k_0T$), and using Eqs. (1) and (2) and the symmetry of the SL in the plane of the layers, the following expressions can be obtained for the transverse ($\tau_\perp=\tau_x=\tau_y$) and longitudinal ($\tau_\parallel=\tau_z$) components of the relaxation time tensor:

$$\frac{1}{\tau_\perp(E)}=\frac{4m^*V}{\pi d\hbar^3}\int_0^1 \frac{W(2k_\perp x)}{\sqrt{1-x^2}}x^2dx, \tag{3}$$

$$\frac{1}{\tau_\parallel(E)}=\frac{2m^*V}{\pi d\hbar^3}\int_0^1 \frac{W(2k_\perp x)}{\sqrt{1-x^2}}dx, \tag{4}$$

where $E=\hbar^2k_\perp^2/2m^*$, $k_\perp=\sqrt{k_x^2+k_y^2}$, $V=SL$ is the volume of a SL of thickness $L=Nd$, and N is the number of periods in the SL. The function W is related to the total probability $P_{\mathbf{kk}'}$ of an electron making a transition from the state $\mathbf{k}$ into the state $\mathbf{k}'$ in a definite type of scattering by the formula

$$P_{\mathbf{kk}'}=\frac{2\pi}{\hbar}W(|\mathbf{k}'_\perp-\mathbf{k}_\perp|)\,\delta[E(\mathbf{k}')-E(\mathbf{k})], \tag{5}$$

where $W(|\mathbf{k}'-\mathbf{k}|)=|U_{\mathbf{k}'\mathbf{k}}|^2$, and $U_{\mathbf{k}'\mathbf{k}}=U(|\mathbf{k}'-\mathbf{k}|)$ is the matrix element of the scattering operator $\hat{U}(\mathbf{r})$ in the basis of Bloch-type envelope wave functions $\tilde{\psi}_{\mathbf{k}}(\mathbf{r})$:

$$U(\mathbf{k}'-\mathbf{k})=\frac{1}{V}\int \tilde{\psi}^*_{\mathbf{k}'}\hat{U}(\mathbf{r})\tilde{\psi}_k d\mathbf{r}, \tag{6}$$

$$\tilde{\psi}_{\mathbf{k}}(r)=\frac{1}{\sqrt{S}}\exp[i(k_x x+k_y y)]\psi_q(z). \tag{7}$$

For the SL considered here, a Bloch sum over wave functions of the ground state of infinitely deep QWs was taken as the function describing electron motion along the axis of the SL:

$$\psi_q(z)=\frac{1}{\sqrt{L}}\exp(iqz)u_q(z)$$

$$=\frac{1}{\sqrt{N}}\sum_{n=-N/2}^{N/2}\exp(iqdn)\varphi(z-dn), \tag{8}$$

$$\varphi_q(z)=\sqrt{\frac{2}{a}}\cos\left(\frac{\pi}{a}z\right)\left[-\frac{a}{2}<z<\frac{a}{2}\right],$$

$$\varphi_q(z)=0\left[z<-\frac{a}{2},\quad \frac{a}{2}<z\right]. \tag{9}$$

It should be noted that using the approximations described above, according to Eq. (5), the scattering probability does not depend on the longitudinal wave vector.

1.1. Scattering by impurity atoms and neutral atoms

For scattering by impurity atoms and neutral atoms the scattering operator has the form

$$\hat{U}(\mathbf{r})=U(\mathbf{r})=\sum_{\alpha}U(\mathbf{r}-\mathbf{R}_\alpha), \tag{10}$$

where $U(\mathbf{r}-\mathbf{R}_\alpha)$ is the interaction energy of an electron with an ion or neutral atom located at the point $\mathbf{R}_\alpha$. Using Eqs. (6)–(10), we have

$$U(\mathbf{k})=\sum_{\alpha}U_\alpha(\mathbf{k}), \tag{11}$$

where

$$U_\alpha(\mathbf{k})=\frac{1}{V}\sum_{n}\exp(-i\mathbf{k}_n\cdot\mathbf{R}_\alpha)S_n(q)U(\mathbf{k}_n), \tag{12}$$

$$S_n(q)=\frac{1}{q}\int_{-d/2}^{d/2}\exp(i2\pi nz/d)u_q^*(z)u_0(z)dz, \tag{13}$$

$$U(\mathbf{k})=U(k_\perp,q)=\int_V\exp(-i\mathbf{k}\cdot\mathbf{r})U(\mathbf{r})d\mathbf{r}, \tag{14}$$

$$\mathbf{k}_n=(\mathbf{k}_\perp,q+2\pi n/d).$$

Using a uniform distribution of scattering centers over the volume of the SL or the volume of the QW, we obtain the following expression for the desired function $W(k_\perp)$ with the formulas presented above:

$$W(k_\perp)=\frac{N_i d}{\pi V}\int_0^\infty|S(q)|^2|U(k_\perp,q)|^2dq, \tag{15}$$

where N_i is the density of scattering centers and

$$S(q)=\frac{\pi^2\sin(aq/2)}{(aq/2)[\pi^2-(ad/2)^2]}. \tag{16}$$

For scattering by a screened Coulomb ion potential, we have in the approximation of isotropic and uniform permittivity

$$|U(k_\perp,q)|^2=\frac{e^4}{\varepsilon_0^2\varepsilon_s^2(k_\perp^2+q^2+\alpha_s^2)^2}, \tag{17}$$

where ε_s is the static permittivity, and α_s^{-1} is the Debye screening length for the static field. For a SL in the two-dimensional electron gas approximation

$$\alpha_s^2=\alpha_{s0}^2\frac{N_c}{n+N_c}, \tag{18}$$

where $\alpha_{s0}^{-1}=(ne^2/\varepsilon_0\varepsilon_s k_0T)^{-1/2}$ is the Debye screening length of a bulk semiconductor for a nondegenerate electron gas, n is the electron density in the SL, and $N_c=m^*k_0T/\pi d\hbar^2$ is the effective density of states of the two-dimensional electron gas in the bottom miniband. For elastic scattering by neutral atoms, according to the theory of Ref. 15, extended to the case of a two-dimensional electron gas,

$$|U(k_\perp,q)|^2=\frac{30\pi r_0\hbar^4}{m^*}(k_\perp^2+q^2+\alpha_s^2)^{-1/2}, \tag{19}$$

where r_0 is the effective Bohr radius.

1.2. Scattering by phonons

As is well known, the operator representing the interaction of an electron with phonons for a separate allowed branch in a bulk semiconductor can be written as

$$\hat{U}(\mathbf{r})=\hat{U}^+(\mathbf{r})+\hat{U}^-(\mathbf{r}), \tag{20}$$

where

$$\hat{U}^\pm(\mathbf{r})=\sum_{\mathbf{Q}}A^\pm(\mathbf{Q})e^{\mp i\mathbf{Q}\cdot\mathbf{r}} \tag{21}$$

are operators corresponding to creation and annihilation of phonons during scattering and $\mathbf{Q}$ is the phonon wave vector. Substituting the expression (21) into Eq. (12) and equating $\mathbf{R}_\alpha$ to zero, we obtain

$$W(k_\perp)=W^+(k_\perp)+W^-(k_\perp), \tag{22}$$

where

$$W^\pm(k_\perp)=|U^\pm(\mathbf{k})|^2=\left(\frac{d}{\pi}\right)^2\left|\int_0^\infty S(q)A^\pm(k_\perp,q)dq\right|^2. \tag{23}$$

1.2.1. Acoustic phonons

For scattering of electrons by longitudinal acoustic phonons in a volume-type SL (Ref. 14), we use in the deformation potential approximation for the average value of the operators $A^{\pm}$ the standard formula

$$A^{\pm}(\mathbf{Q}) = \mp iQD\sqrt{\frac{\hbar}{2\rho V\omega_L(\mathbf{Q})}\left[N_L(\mathbf{Q})+\frac{1}{2}\pm\frac{1}{2}\right]}, \quad (24)$$

where D is the deformation-potential constant, ρ is the density of the crystal, $\omega_L = v_L Q$ and v_L are, respectively, the frequency and velocity of longitudinal long-wavelength acoustic phonons, and $N_L(\mathbf{Q})$ is the thermodynamically equilibrium number of phonons, described by the Bose–Einstein distribution function. For $\hbar\omega_L \ll k_0 T$

$$A^{\pm}(\mathbf{Q}) = A = i\frac{D}{v_L}\sqrt{\frac{k_0 T}{2\rho V}}. \quad (25)$$

Using Eqs. (22), (23), and (25), we obtain for the desired function

$$W(k_\perp) = 2\left(\frac{Ad}{\pi}\right)^2\left|\int_0^\infty S(q)dq\right|^2 = 4\frac{k_0 T}{\rho V}\left(\frac{Dd}{v_L a}\right)^2. \quad (26)$$

1.2.2. Optical polar phonons

For scattering of electrons by longitudinal optical polar phonons of the volume type[12,13] we use for the average value of the operators $A^{\pm}(\mathbf{Q})$ the well-known formula corresponding to a screened optical potential and the high-temperature approximation

$$A^{\pm}(\mathbf{Q}) = \mp A(\mathbf{Q}) = \mp i\frac{eP_L}{\omega_L^{\mathrm{PO}}}\sqrt{\frac{k_0 T}{2\varepsilon_0 V}}\frac{Q}{Q^2+\alpha_\infty^2}, \quad (27)$$

where ω_L^{PO} is the frequency of long-wavelength longitudinal optical polar phonons, p_L is the oscillator strength, and α_∞^{-1} is the Debye screening length for the high-frequency field, which differs from α_s^{-1} by the fact that in Eq. (18) the static permittivity ε_s is replaced by the high-frequency permittivity ε_∞. As is well known, the following relation holds for III–V semiconductors:

$$\frac{P_L^2}{\omega_L^2} = \frac{1}{\varepsilon^*} = \frac{1}{\varepsilon_\infty} - \frac{1}{\varepsilon_s}. \quad (28)$$

Using Eqs. (23) and (27), we obtain the following expression for the desired function:

$$W(k_\perp) = 2\left(\frac{d}{\pi}\right)^2\left|\int_0^\infty S(q)A(k_\perp, q)dq\right|^2. \quad (29)$$

1.3. Electron momentum relaxation time

Using the formulas obtained for the function $W(k_\perp)$, we represent the components of the tensor of the reciprocal of the momentum relaxation time for individual scattering mechanisms in the form

$$\frac{1}{\tau_i(E)} = \frac{1}{\tau_0(E)}G_i(E), \quad (30)$$

where

$$\tau_0(E) = (ak_\perp)\tau(E), \quad (31)$$

$\tau(E)$ is the relaxation time in a bulk semiconductor, and $G_i(E)$ is a dimensionless function, whose parameters, in general, are the QW width, the SL period, and the free charge-carrier density.

1.3.1. Impurity ions

For scattering by impurity ions, using Eqs. (3), (4), and (15)–(17), we obtain for the functions $G_i(E)$ the integral expressions

$$G_\perp(E) = \frac{\pi^4}{2F(\eta)}\int_0^\infty \frac{dx\sin^2 x}{A(x,\gamma,\beta_s)(\gamma^2x^2+\beta_s^2)^{1/2}}, \quad (32)$$

$$G_\parallel(E) = \frac{\pi^4}{2F(\eta)}\int_0^\infty \frac{dx(2\gamma^2x^2+2\beta_s^2+1)\sin^2 x}{A(x,\gamma,\beta_s)(\gamma^2x^2+\beta_s^2)^{3/2}}, \quad (33)$$

$$A(x,\gamma,\beta_s) = x^2(\pi^2-x^2)(\gamma^2x^2+\beta_s^2+1)^{3/2},$$

where $\gamma = 1/ak_\perp$, $\beta_s = \alpha_s/2k_\perp$, and the function $F(\eta) = \ln(1+\eta) - \eta/(1+\eta)$, which depends on the parameter $\eta = (2k_\perp/\alpha_{s0})^2$, is known from the theory of Coulomb scattering in bulk semiconductors.

1.3.2. Neutral impurity atoms

For scattering by neutral impurity atoms, using Eqs. (3), (4), (15), (16), and (19), we can write the functions $G_i(E)$ as a pair of integrals

$$G_\perp(E) = \frac{4}{\pi}\int_0^1 \frac{x^2}{\sqrt{1-x^2}}g(E,x)dx,$$

$$G_\parallel(E) = \frac{2}{\pi}\int_0^1 \frac{1}{\sqrt{1-x^2}}g(E,x)dx, \quad (34)$$

where

$$g(E,x) = \frac{3\pi^4}{2}\int_0^\infty \frac{dy\sin^2 y}{y^2(\pi^2-y^2)^2(\gamma^2y^2+x^2+\beta_s^2)^{1/2}}. \quad (35)$$

1.3.3. Acoustic phonons

For scattering by acoustic phonons, using Eqs. (3), (4), and (26), we found the functions $G_i(E)$ to be constant:

$$G_\perp(E) = G_\parallel(E) = 4\pi\frac{d}{a}. \quad (36)$$

1.3.4. Polar optical phonons

Using Eqs. (3), (4), and (29), we can represent the functions $G_i(E)$ for scattering by polar optical phonons in the form of Eqs. (34), where

$$g(E,x) = \pi^3\left(\frac{d}{a}\right)\left|\int_0^\infty \frac{dy\sin y(\gamma^2y^2+x^2)^{1/2}}{y(\pi^2-y^2)(\gamma^2y^2+x^2+\beta_\infty^2)}\right|^2,$$

$$\beta_\infty = \alpha_\infty/2k_\perp. \quad (37)$$

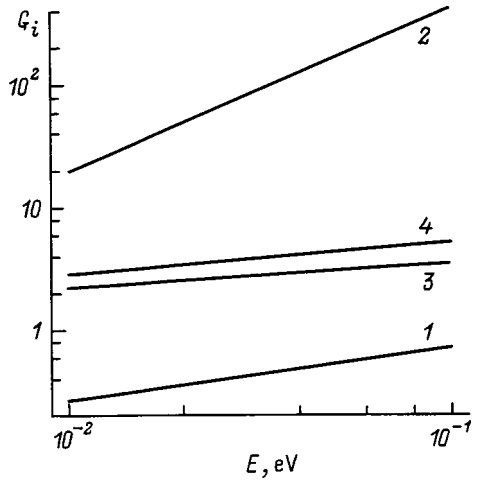

FIG. 1. Dimensionless function G_i versus the energy for $T=100$ K. $a=b=5$ nm, $n=10^{15}$ cm^{-3} for scattering: *1, 2*—by impurity ions and *3, 4*—by neutral atoms. *1, 3*—Transverse component of the tensor, *2, 4*—longitudinal component of the tensor.

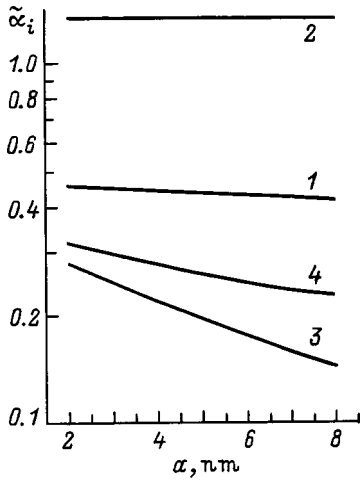

FIG. 2. $\tilde{\alpha}_i$ versus the quantum well width at $T=100$ K for $b=5$ nm and $n=10^{15}$ cm^{-3}. The numbers on the curves have the same meaning as in Fig. 1.

2. MOBILITY TENSOR

Using Eqs. (1) and (2), we can represent the diagonal components of the mobility tensor $\hat{\mu}$, written in principal axes, in the standard form

$$\mu_\perp = \frac{e\langle \tau_\perp \rangle}{m^*}, \quad \mu_\parallel = \frac{e\langle \tau_\parallel \rangle}{\langle m_\parallel \rangle}, \tag{38}$$

where $\langle \tau_\perp \rangle$ and $\langle \tau_\parallel \rangle$ are the energy-averaged transverse and longitudinal relaxation times, $\langle m_\parallel \rangle$ is the average longitudinal effective mass,

$$\langle \tau_\perp \rangle = \frac{1}{n}\int_0^\infty \left(-\frac{\partial f_0}{\partial E}\right)\tau_\perp(E)\rho_c(E)E\,dE, \tag{39}$$

$$\langle \tau_\parallel \rangle = [1-\exp(-n/N_c)]^{-1}\int_0^\infty \left(-\frac{\partial f_0}{\partial E}\right)\tau_\parallel(E)\,dE, \tag{40}$$

$$\frac{1}{\langle m_\parallel \rangle} = \frac{\Delta\rho_c}{4m_\parallel n}[1-\exp(-n/N_c)]. \tag{41}$$

In Eqs. (39)–(41), $\rho_c = m^*/\pi d\hbar^2$ is the density of states in the bottom conduction miniband in the two-dimensional electron gas approximation, $N_c = k_0 T\rho_c$ is the effective density of states in the bottom miniband, and $m_\parallel = 2\hbar^2/\Delta^2 d^2$ is the longitudinal effective mass at the miniband bottom. For a nondegenerate electron gas, which corresponds to satisfying the condition $n \ll N_c$, in the power-law approximation for the relaxation time as a function of the energy and temperature

$$\tau_i(E) = \tau_{0i}(k_0T)^{\beta_i}E^{\alpha_i+1/2}, \tag{42}$$

the expressions for the average relaxation times (39) and (40) and for the average longitudinal effective mass (41) assume the form

$$\langle \tau_i \rangle = \tau_{0i}(k_0T)^{\alpha_i+\beta_i+1/2}\Gamma(\alpha_i+\delta_i), \tag{43}$$

$$\frac{1}{\langle m_\parallel \rangle} = \frac{\Delta}{4k_0T}\frac{1}{m_\parallel} \sim T^{-1}, \tag{44}$$

where $\Gamma(n)$ is the gamma function, $\delta_\perp = 5/2$, and $\delta_\parallel = 3/2$. Using Eqs. (38), (43), and (44), we obtain for the temperature dependence of the mobility

$$\mu_i \sim T^{\gamma_i}, \tag{45}$$

where $\gamma_i = \alpha_i + \beta_i + \chi_i$, $\chi_i = 1/2$, and $\chi_\parallel = -1/2$.

For a degenerate electron gas with $n \gg N_c$ we obtain, instead of Eqs. (43) and (44), formulas that depend on the reduced Fermi level ξ, or electron density $n = \rho_c\xi$,

$$\langle \tau_i \rangle = \tau_i(\xi) \sim T^{\beta_i}\xi^{\alpha_i+1/2} \sim T^{\beta_i}n^{\alpha_i+1/2}, \tag{46}$$

$$\frac{1}{\langle m_\parallel \rangle} = \frac{\Delta}{4\xi}\frac{1}{m_\parallel} \sim n^{-1}. \tag{47}$$

According to Eqs. (46) and (47), we obtain the following expression for the temperature and electron density dependence of the mobility:

$$\mu_i \sim T^{\beta_i}n^{\alpha_i+\chi_i}. \tag{48}$$

3. NUMERICAL ANALYSIS FOR SCATTERING BY IMPURITY IONS AND NEUTRAL ATOMS

Since analytic expressions could not be obtained for the formulas describing scattering by impurity ions and neutral atoms, we analyzed numerically for these scattering mechanisms the energy dependence of the relaxation time and temperature dependence of the mobility. The analysis was made in the temperature range $50 < T < 300$ K for a composite SL with the parameters $a=b=5$ nm, $m^*=0.1m_0$, potential barrier height $V_{SL}=0.4$ eV, $\varepsilon_s=\varepsilon_\infty=10$, and $n=10^{15}$ cm^{-3}, which are characteristic of SLs used in infrared-range photodetectors.[1,2] According to numerical calculations,[16] the energy spectrum of this SL in the QW has two minibands of width $\Delta_1=1.6$ meV and $\Delta_2=18.5$ meV, separated in energy by 210 meV. These data substantiate the validity of the approximation, used in the calculations presented above, of a two-dimensional electron gas and a single miniband approximation for the SL.

The analysis associated with the numerical calculation of the functions $G_i(E)$ showed the following.

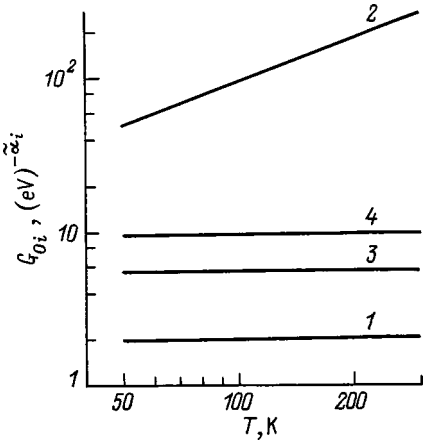

FIG. 3. G_{0i} versus the temperature for $a=b=5$ nm and $n=10^{15}$ cm^{-3}. The numbers on the curves have the same meaning as in Fig. 1.

1. For scattering by impurity ions and neutral atoms the functions $G_i(E)$ show a strong energy dependence, similar to a power-law dependence (see Fig. 1). In other words,

$$G_i(E) \approx G_{0i}(T) E^{\tilde{\alpha}_i}. \tag{49}$$

2. The exponent $\tilde{\alpha}_i$ in this dependence is essentially independent of temperature and the electron density in the range $n < 10^{16}$ cm^{-3}. For scattering by neutral atoms it depends appreciably on the QW width a (see Fig. 2).

3. As the temperature increases, the quantity $G_{0\parallel}$ for scattering by impurity ions (Fig. 3, curve *2*) increases almost linearly, while $G_{0\perp}$ depends very weakly on temperature (Fig. 3, curve *1*). For scattering by neutral atoms this dependence is also very weak for both components (see Fig. 3, curve *3* and *4*).

4. For scattering by impurity ions, $G_{0\parallel}$ was found to be several orders of magnitude greater than G_{0i}, and it increases with decreasing electron density (Fig. 4, curve *2*). For $G_{0\perp}$ this dependence is weak. For scattering by neutral atoms, $G_{0\parallel}$ and $G_{0\perp}$ are of the same order of magnitude and are virtually independent of electron density (see Fig. 4).

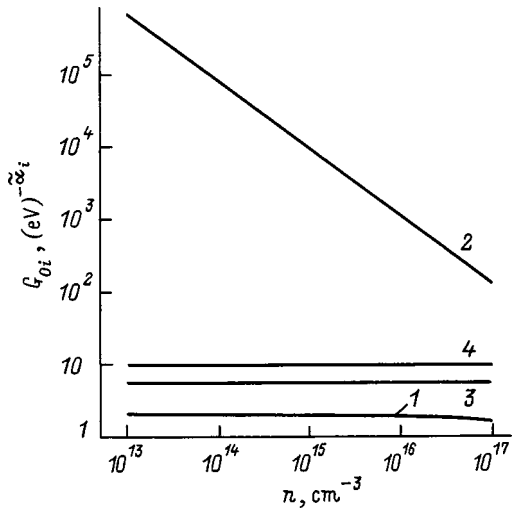

FIG. 4. G_{0i} versus the electron density at $T=100$ K for $a=b=5$ nm. The numbers on the curves have the same meaning as in Fig. 1.

TABLE I. Exponents in the energy and temperature dependences of the momentum relaxation time and mobility of electrons in a superlattice.

Scattering mechanism	$\alpha_0+\beta_0$	$\tilde{\alpha}_i$		$\tilde{\beta}_i$		γ	
		$i=\perp$	$i=\parallel$	$i=\perp$	$i=\parallel$	$i=\perp$	$i=\parallel$
Impurity ions	1.5	0.4	1.4	0	0.9	1.6	−1.3
Neutral atoms	0	0.2	0.3	0	0	0.3	−0.8
Acoustic phonons	−1.5	0	0	0	0	−1.0	−2.0

Note: The exponents refer to the transverse ($i=\perp$) and longitudinal ($i=\parallel$) components of the tensor.

On the basis of the above analysis, the dependence of the functions G_i on temperature and energy can be represented in the form

$$G_i \sim T^{\tilde{\beta}_i} E^{\tilde{\alpha}_i}, \tag{50}$$

Then, using Eqs. (31) and (45), the temperature-dependence of the mobility for a nondegenerate electron gas assumes the form (45), where $\alpha_i=\alpha_0-\tilde{\alpha}_i$, $\beta_i=\beta_0-\tilde{\beta}_i$, and α_0 and β_0 are the exponents in the energy and temperature dependences of the relaxation time in a bulk semiconductor. The temperature-average values of $\tilde{\alpha}_i$, $\tilde{\beta}_i$, and γ_i, calculated with the above-indicated parameters of the SL and corresponding to the three basic scattering mechanisms, are presented in Table I. It follows from the table that a strong change in the temperature-dependence of the mobility in the SL under study, in contrast to its constituent bulk semiconductors, should be expected for the longitudinal component. In the extrinsic conductivity range, determined by scattering by impurity ions and neutral atoms, this dependence for the longitudinal mobility is of a qualitatively different character than in bulk semiconductors, since it should decrease with increasing temperature.

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Translated by M. E. Alferieff