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Momentum relaxation of 2D electron gas due to near-surface acoustic phonon scattering☆

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Abstract

This study addresses the effect of proximity of a quantum well to a stress-free surface of the semiconductor heterostructure on the momentum relaxation rate of two-dimensional electrons interacting with acoustic phonons via piezoelectric and deformation potentials. The results obtained demonstrate that for narrow quantum wells placed close to the surface the relaxation rate at low temperatures (Bloch–Grüneisen regime) is changed considerably in comparison with that of a two-dimensional electron gas placed in a bulk of semiconductor. For the temperatures where the piezoelectric potential interaction dominates over the deformation potential interaction, the near-surface relaxation rate is enhanced in the case of a semiconductor–vacuum system and is suppressed in the case of the surface covered by a thin metal film. The temperature dependence of the near-surface momentum relaxation rate is found to be T^{α} for values of T far below the Bloch–Grüneisen temperature. For a semiconductor–vacuum system, $\alpha = 3$ and 5 for piezoelectric and deformation potential scattering, respectively; for a semiconductor–metal system, $\alpha = 5$ for both mechanisms. It is predicted that screening changes the temperature dependences of momentum relaxation rates: for a semiconductor–vacuum system, $\alpha = 5$ and 7 for piezoelectric and deformation potential scattering, respectively. Screening does not change α in the case of metal–semiconductor system.

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1. Introduction

In recent years, the influence of the surface of a semiconductor heterostructure with an acoustic phonon scattering of two-dimensional electron gas (2DEG) attracted substantial interest [1–7]. The effect to be examined in this account is caused by the difference in the acoustic phonon modes in a bounded medium and the conventional

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three-dimensional (3D) phonon modes; it also depends on the mechanical boundary conditions at the surfaces and on the geometry of the system. Kinetic phenomena associated with a 2DEG in free-standing structures (see e.g. Ref. [2] and references therein), in slabs [3], and in semi-infinite semiconductors [1,4–7] have been studied previously. It was shown [5–7] that for a 2DEG localized at a finite distance from the surface, the modification of the acoustic displacement field may substantially change the relaxation processes compared with that for 2DEG placed in the bulk of semiconductor. For electron–acoustic phonon interaction via the deformation potential (DP), a large influence on energy and momentum relaxation rates occurs at low temperatures when the scattering processes are inelastic (Bloch–Grüneisen regime). This regime corresponds to temperatures T which are less than or comparable to the characteristic temperature $T_0 = 2sp_F/k_B$ where s, depending on the phonon mode involved, is the longitudinal or transversal sound velocity, p_F the electron Fermi momentum, and k_B Boltzmann's constant. In the high-temperature range, $T \gg T_0$, the scattering processes are quasi-elastic and, accordingly, we can consider the equipartition energy distribution of the phonons. In this case, the sensitivity of the scattering to the phonon mode structure becomes weaker; indeed, the presence of the surface results in a small change (approximately 10%) of the relaxation rates, and the conventional bulk phonon results for the mobility and the energy relaxation rates of 2DEG [8] are appropriate. A small change of the same order in mobility has been reported earlier [9] for a 2DEG located at a Si/SiO_2 interface using a soft-medium approximation for SiO_2.

In compound semiconductors without inversion symmetry, electrons interact with acoustic phonons through both deformation and piezoelectric potentials. The studies [10,11] of piezoelectric acoustic (PA) scattering, which is governed by bulk phonons, demonstrate that this scattering may dominate over that of the DP interaction at low temperatures. For GaAs-based heterostructures, the transition from DP- to PA-coupled phonon scattering takes place for temperatures close to the transition temperature T_0; a typical quoted value of T_0 is 6 K [12]. Thus, to study the influence of the surface on the electron kinetics in real semiconductor materials, one also has to take into account the PA scattering. Both piezoelectric and deformation potential interactions of surface acoustic phonons (Rayleigh waves) with 2DEG have been considered in Ref. [4]; the corresponding transport relaxation rates for 2DEG were calculated in Ref. [5]. To the best of our knowledge, the total near-surface PA scattering of a 2DEG interacting with the full set of phonon modes in semi-bounded systems has not been carried out to this date. As stated previously, the subject of this paper is the study of the effect of proximity of a 2DEG to a surface of a heterostructure on the transport relaxation rates due to both the PA and DP interactions. The influence of a surface on PA scattering is not restricted to the effects associated with the modification of phonons; indeed, it also depends on the boundary conditions for the piezoelectric potential. We consider a semi-bounded heterostructure with a stress-free surface and we use two types of electrical boundary conditions; one corresponds to the semiconductor–vacuum system and the other to a semiconductor covered by a thin metal film.

2. Model and basic equations

We shall consider a semiconductor heterostructure which occupies the half-space $z > 0$ with a stress-free plane boundary at $z = 0$. At the distance z_0 from the surface, there is a 2D electron channel formed by the electrons in a rectangular quantum well (QW) of width d, see Fig. 1. For the sake of simplicity, we assume the same elastic and piezoelectric properties, densities, and dielectric constants for all of the layers of the heterostructure. The potential outside the QW is taken to be infinite so that the proximity of the surface $z = 0$ to the QW does not disturb the stationary electron states, and the influence of the surface on the electron subsystem appears only as a result of the scattering with acoustic phonons modified by the surface. We assume that electrons occupy the lowest subband, and that the wave function for transverse movement is $\psi(z) = (2/d)^{1/2}\cos(\pi(z - z_0)/d)$ for $|z - z_0| \leqslant d/2$ and $\psi(z) = 0$ outside the QW. The

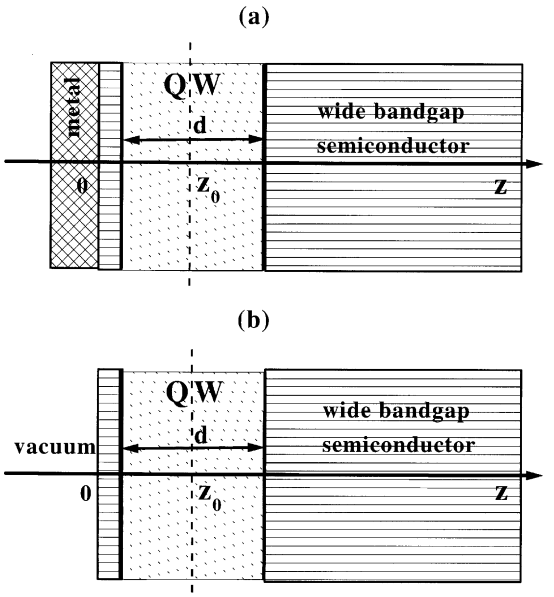

Fig. 1. Schematic representation of the structure under consideration. Two-dimensional electron channel of width d at the distance z_0 from the surface of: (a) semiconductor–metal boundary; (b) semiconductor–vacuum boundary.

momentum relaxation rate, ν_m, is introduced by the balance equation

$$\frac{2}{n_s L^2}\sum_{p}\frac{\boldsymbol{p}}{m^*}J_{e\text{-}ph}(\boldsymbol{p}) = -\nu_m \boldsymbol{v}_d. \tag{1}$$

Here $\boldsymbol{p}=(p_x,p_y)$ is the 2D momentum of the electrons, n_s is the electron sheet density, L^2 is the normalization area in the xy-plane, m^* and $\boldsymbol{v}_d$ are the electron effective mass and the drift velocity, respectively. The integral of electron–phonon collisions, $J_{e\text{-}ph}(\boldsymbol{p})$, is related to the transition probability and to the nonequilibrium electron distribution function in usual way. Using a shifted Fermi distribution function, in the case of a small deviation from thermodynamic equilibrium we get

$$\nu_m = \frac{1}{n_s L^2}\sum_{pp'}\frac{[(\boldsymbol{p}-\boldsymbol{p}')\cdot\boldsymbol{v}_d]^2}{k_B T m^* |\boldsymbol{v}_d|^2} W(\boldsymbol{p},\boldsymbol{p}') f_\varepsilon(1-f_{\varepsilon'}). \tag{2}$$

Here $W(\boldsymbol{p},\boldsymbol{p}')$ is the transition probability for a transition from the electron's initial state with 2D momentum $\boldsymbol{p}$ to the final state with momentum $\boldsymbol{p}'$ within the first subband due to the interaction of an electron with acoustic phonon modes, and f_ε is the equilibrium Fermi distribution function; the energy spectrum is taken to be parabolic: $\varepsilon = |\boldsymbol{p}|^2/2m^*$. The electrons are assumed to be degenerate, i.e $T \ll \pi n_s \hbar^2/(m^* k_B)$. The drift velocity is expressed in terms of the rate ν_m as $\boldsymbol{v}_d = (e\nu_m^{-1}/m^*)\boldsymbol{E}$, where $\boldsymbol{E}$ is the strength of the applied electric field. If one uses for the nonequilibrium distribution function the ordinary expansion, $f_\varepsilon - \tau(p)(\boldsymbol{p}\cdot\boldsymbol{E})\mathrm{d}f_\varepsilon/\mathrm{d}\varepsilon$, with the electron momentum relaxation time τ instead of the shifted Fermi distribution function model, then for degenerate electrons, the average relaxation time, ν_m^{-1}, in Eq. (2) equals $\tau(p_F)$. We take into account inelastic scattering in Eq. (2).

The interaction Hamiltonian is given by

$$H_{int} = e\phi + D\,\mathrm{div}\,\boldsymbol{u}, \tag{3}$$

where D is the deformation potential constant and ϕ is the piezoelectric potential. The displacement vectors, $\boldsymbol{u}$, are found from the elastic wave equation in the isotropic continuum approximation:

$$\frac{\partial^2}{\partial t^2}\boldsymbol{u} = s_t^2\nabla^2\boldsymbol{u} + (s_l^2 - s_t^2)\nabla(\nabla\cdot\boldsymbol{u}), \tag{4}$$

where s_l and s_t are the velocities of the longitudinal and transversal waves, respectively. For cubic semiconductors, these velocities are taken to be equal to the appropriate average quantities [13]. The stress-free boundary conditions, namely $\sigma_{xz} = \sigma_{yz} = \sigma_{zz} = 0$ for the stress tensor σ at the surface $z=0$, and the requirement that the solutions to be finite at $z\to\infty$, are imposed. Taking into account the piezoelectric polarization, Poisson's equation has the following form:

$$\nabla^2\phi = \beta\left(\frac{\partial^2 u_x}{\partial y\partial z} + \frac{\partial^2 u_y}{\partial x\partial z} + \frac{\partial^2 u_z}{\partial x\partial y}\right), \quad z \geqslant 0. \tag{5}$$

We assume that the z-axis is oriented along the [0 0 1] direction in a cubic crystal. Accordingly, the tensor of the piezoelectric moduli has only one non-zero component, and $\beta = 8\pi e_{14}/\varepsilon_0$, where ε_0 is the lattice dielectric permittivity, and e_{14} is the component of the tensor which relates the strength of an electric field to the strain tensor.

We shall consider the two types of boundary conditions which are of interest in practice:

open-circuit and short-circuit conditions [14]. The open-circuit condition corresponds to the case when the surface of the semiconductor has an electrically free boundary with a vacuum. For this condition, the electric potential associated with the acoustic waves vanishes in vacuum as the distance from the crystal increases, and the potential and normal components of dielectric displacement are continuous at the surface $z = 0$:

$$\phi(z = -0) = \phi(z = +0),$$

$$\left(\frac{\partial\phi}{\partial z}\right)_{z=-0} = \varepsilon_0\left[\frac{\partial\phi}{\partial z} - \frac{\beta}{2}\left(\frac{\partial u_x}{\partial y} + \frac{\partial u_y}{\partial x}\right)\right]_{z=+0}. \quad (6)$$

The short-circuit condition is applied for a piezoelectric sample covered with a metal film which is thin enough for the stress-free mechanical boundary conditions to be unchanged. In this case, the electric potential obeys the condition

$$\phi(z = 0) = 0. \quad (7)$$

To introduce acoustic phonons in a semi-bounded medium, one has to construct a complete set of the orthonormalized vectors $\boldsymbol{u}(\boldsymbol{r}, t)$ by solving the eigenvalue problem of the elastodynamic equation (4). Then the phonons are determined using the second quantization formalism. The full set of phonon modes in such a semi-infinite medium is given in Ref. [15]. In our calculations, we exploit the more convenient set of modes used in Ref. [7]. In this representation, phonons are characterized by the set of quantum numbers $\omega, \boldsymbol{q}$, and j, where ω is the angular frequency, $\boldsymbol{q} = (q_x, q_y)$ is the in-plane wave vector, and the label j specifies different types of the modes. To describe the PA interaction, we generalized the results of Ref. [7] to the case of an arbitrarily directed wave vector, $\boldsymbol{q}$. The jth mode is expressed as

$$\boldsymbol{u}_j(\boldsymbol{r}, t) = \boldsymbol{u}_j(z)\mathrm{e}^{\mathrm{i}(\boldsymbol{q}\boldsymbol{r} - \omega t)},$$

where the vectors $\boldsymbol{u}_j(z)$ are some linear combinations of the functions $\exp(\pm \mathrm{i}k_v z)$ and $\exp(-\kappa_v z)$; $k_v = (\omega^2/s_v^2 - q^2)^{1/2}$ and $\kappa_v = (q^2 - \omega^2/s_v^2)^{1/2}$ are determined for $\omega > s_v q$ and $\omega < s_v q$, respectively; $v = \mathrm{l}, \mathrm{t}$ corresponds to longitudinal and transverse waves, respectively.

The modes are found in different ways depending on the frequency ranges. In the range $\omega > s_\mathrm{t} q$, the vectors $\boldsymbol{u}_j(z)$ correspond to the different possibilities for waves incident from the bulk upon the surface. For $\omega < s_\mathrm{t} q$, there are no incident waves, and the solution is represented by linear combination of decaying waves only.

In the second quantization representation, the phonon displacement operator is

$$\hat{\boldsymbol{u}} = \sum_{j,\boldsymbol{q}} \int \mathrm{d}\omega \left(\frac{\hbar}{2\rho\omega L^2}\right)^{1/2} \times [\boldsymbol{u}_{j\omega\boldsymbol{q}}(z)\mathrm{e}^{\mathrm{i}(\boldsymbol{q}\boldsymbol{r} - \omega t)} b_{j\omega\boldsymbol{q}} + \mathrm{H.c.}], \quad (8)$$

where $b_{j\omega\boldsymbol{q}}$ is a phonon annihilation operator, ρ represents the density, and $j = \mathrm{l}, \mathrm{th}, \mathrm{tv}, \mathrm{R}$. The first three modes correspond to the following choice of incident waves: longitudinal incident wave (l-mode), transverse horizontal wave polarized in xy-plane (th-mode), and transverse vertical wave polarized in the plane of incidence (tv-mode); the R-mode denotes a Rayleigh wave which has velocity s_R and obeys the dispersion law $\omega = s_\mathrm{R} q$. For each mode, the integration is over the range where the mode exists.

Substituting Eq. (8) in Eqs. (5)–(7) we get the equations and boundary conditions for the phonon induced potentials, $\phi_{j\omega\boldsymbol{q}}(z)$. The interaction Hamiltonian of Eq. (3) is rewritten as the interaction operator

$$\hat{H} = \sum_{j,\boldsymbol{q}} \int \mathrm{d}\omega \left(\frac{\hbar}{2\rho\omega L^2}\right)^{1/2} \left[e\phi_j + D\left(\mathrm{i}\boldsymbol{q}\boldsymbol{u}_j + \frac{\partial u_{jz}}{\partial z}\right)\right] \times \mathrm{e}^{\mathrm{i}(\boldsymbol{q}\boldsymbol{r} - \omega t)} b_j + \mathrm{H.c.} \quad (9)$$

For convenience we drop the subscript $\omega\boldsymbol{q}$ in the notations for ϕ_j, $\boldsymbol{u}_j$, and b_j.

In order to calculate the momentum relaxation rate, one has to find the transition rate $W(\boldsymbol{p}, \boldsymbol{p}')$ using perturbation theory with Eq. (9), and then carry out the summation over the initial and final electron momenta $\boldsymbol{p}$ and $\boldsymbol{p}'$, respectively.

3. Results and discussion

The transition probability due to the interaction with the jth mode is calculated within the Fermi

golden rule approximation:

$$W_j^{\pm}(\boldsymbol{p},\boldsymbol{p}') = \frac{2\pi}{\hbar}\int \mathrm{d}\omega \frac{\hbar}{2\rho\omega L^2}|M_j|^2(N_\omega + \tfrac{1}{2} \pm \tfrac{1}{2}) \times \delta_{\boldsymbol{p}',\boldsymbol{p}\mp\hbar\boldsymbol{q}}\delta(\varepsilon' - \varepsilon \pm \hbar\omega), \quad (10)$$

where the upper (lower) sign corresponds to phonon emission (absorption), $N_\omega = [\exp(\hbar\omega/k_\mathrm{B}T) - 1]^{-1}$ is the equilibrium phonon distribution function, and M_j is the scattering matrix element

$$M_j(\omega,\boldsymbol{q}) = \int_0^\infty \mathrm{d}z\psi(z)^2\left[e\phi_j + D\left(\mathrm{i}\boldsymbol{q}u_j + \frac{\partial u_{jz}}{\partial z}\right)\right]. \quad (11)$$

The l, tv, and R-modes contribute to both the DP and the PA interactions, and the th-mode contributes to PA interaction only. For each mode, the electron–acoustic–phonon interaction is determined by the square of the absolute value of the scattering matrix element given by Eq. (11). For a 2DEG placed in a bulk of semiconductor, collisions with ordinary 3D phonons increase the relaxation of electron momentum. In this case, the DP and PA mechanisms give additive contributions to the transition rate, so that Matthiessen's rule for the mobility is valid. In contrast to the case of bulk phonons, the partial contributions of DP and PA interactions to the matrix elements, M_j^{DP} and M_j^{PA}, for the scattering on phonons in semi-bounded medium are not out of phase; thus, an interference between these two interaction mechanisms may take place. For the case of scattering on Rayleigh waves, such a possibility has been shown in Ref. [4]. However, in the case of interest, when the surface of cubic crystal is spanned by two lattice axes, one can use the isotropic (in the xy-plane) model. In the framework of this model, the DP and PA mechanisms contribute to the transition rate additively. Indeed, the quantities M_j^{DP} depend on the length of the phonon wave vector q, while the direction-dependent values M_j^{PA} are proportional to the product, q_xq_y, (for $j = \mathrm{l,tv,R}$) or the difference, $q_x^2 - q_y^2$ (for $j = \mathrm{th}$). Upon averaging of $|M_j|^2$ over an azimuthal angle, the interference terms which are linear in q_xq_y vanish and the quantities $(q_xq_y)^2$ and $(q_x^2 - q_y^2)^2$ are replaced by $q^4/8$ and $q^4/2$, respectively. Thus, in the case of the isotropic model, $|M_j|^2 = |M_j^{\mathrm{PA}}|^2 + |M_j^{\mathrm{DP}}|^2$.

Upon performing the integration in Eq. (10) we obtain

$$W_j^{\pm}(\boldsymbol{p},\boldsymbol{p}') = \frac{\pi}{\hbar\rho L^2\omega}|M_j(\omega,q)|^2(N_\omega + \tfrac{1}{2} \pm \tfrac{1}{2}) \quad (12)$$

for the extended modes ($j = \mathrm{l,tv,th}$) and

$$W_j^{\pm}(\boldsymbol{p},\boldsymbol{p}') = \frac{\pi}{\hbar\rho L^2\omega}|M_j(\omega,q)|^2 \times (N_\omega + \tfrac{1}{2} \pm \tfrac{1}{2})\delta(\omega - s_\mathrm{R}q) \quad (13)$$

for the localized modes ($j = \mathrm{R}$). Here and for the remainder of this paper, ω and q are expressed through the energy and momentum transfer of the electrons

$$\hbar\omega = |\varepsilon - \varepsilon'|, \qquad \hbar q = |\boldsymbol{p} - \boldsymbol{p}'| \quad (14)$$

in accordance with the energy and momentum conservation laws. Assuming that $|\varepsilon - \varepsilon'| \ll \varepsilon_\mathrm{F}$ and that the electron velocity is much larger than the speed of sound ($p_\mathrm{F}/m^* \gg s_\mathrm{l}$), we put $p = p' = p_\mathrm{F}$ in relation to q so that $q = 2k_\mathrm{F}\sin\theta$. Here 2θ is the angle between the electron's initial and final momenta, $\boldsymbol{p}$ and $\boldsymbol{p}'$, respectively.

To calculate the momentum relaxation rate, we substitute in Eq. (2) the transition probability, $W(\boldsymbol{p},\boldsymbol{p}')$, as a sum of the partial probabilities of Eqs. (12) and (13). Due to their dependences on the variables of Eq. (14) only, further derivation of the total relaxation rate may be carried out analogously that for DP scattering [16]; the procedure yields the following final form:

$$\nu_\mathrm{m} = \frac{m^*}{8\pi\hbar\rho k_\mathrm{F}^3 k_\mathrm{B}T}\sum_j \int_0^\infty \mathrm{d}\omega \int_0^{2k_\mathrm{F}} \mathrm{d}q \times \frac{q^2|M_j(\omega,q)|^2}{\sqrt{1 - (q/2k_\mathrm{F})^2}\sinh^2(\hbar\omega/2k_\mathrm{B}T)}. \quad (15)$$

The integration domain is divided into three regions each of which is determined by the existence conditions of the corresponding phonon modes; see Fig. 2. In range I ($\omega > s_\mathrm{l}q$), the normal components of wave vectors, $k_{\mathrm{l,t}} = (\omega^2/s_{\mathrm{l,t}}^2 - q^2)^{1/2}$, of bulk longitudinal and transversal waves which form the phonon modes, are real. In this range, the values of

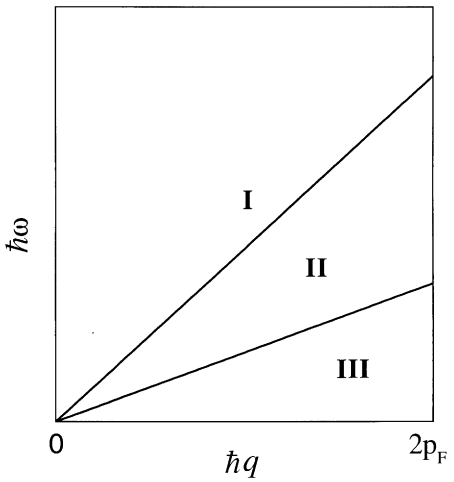

Fig. 2. Domain of integration over the energy transfer, $\hbar\omega = |\varepsilon - \varepsilon'|$, and the momentum transfer, $\hbar q = |\boldsymbol{p} - \boldsymbol{p}'|$, $0 \leqslant \omega < \infty$, $0 \leqslant q \leqslant 2k_F$. The ranges I–III specify the domain of existence of the corresponding phonon modes as described in the text.

M_j are non-zero for $j = \mathrm{l}, \mathrm{tv}, \mathrm{th}$. The interval II ($s_\mathrm{l} q < \omega < s_\mathrm{t} q$) corresponds to real k_t and imaginary k_l. Here, $M_j \neq 0$ for $j = \mathrm{tv}, \mathrm{th}$. In range III ($\omega < s_\mathrm{t} q$) both k_l and k_t are imaginary, and the scattering rate is solely defined by the confined phonons of R-mode. Here the integration is along the line $\omega = s_\mathrm{R} q$, and the contribution of range III is described by a one-dimensional integral. For the particular case of DP interaction, the total matrix elements for the modes involved in each frequency range, coincide with those that were obtained in Ref. [1] using Green functions from the theory of elasticity.

To obtain accurate results, the integrals in Eq. (15) have to be calculated numerically. Let us first discuss the character of the temperature dependence of the momentum relaxation rates for the different types of electrical boundary conditions. In the low-temperature case, when $T \ll T_0$, one can use the small-angle scattering approach. For a narrow QW placed close to the surface and for $T \ll T_0$, the phonon wavelengths are much larger than the width of the well, d, and the distance z_0. The first non-zero term in the expansion of M_j^{PA} over $kz_0 \ll 1$, where $k = k_\mathrm{l}, k_\mathrm{t}, q$ is $(kz_0)^0$ in the case of boundary condition of Eq. (6) and is $(kz_0)^1$ for condition of Eq. (7). The corresponding expansion of M_j^{DP} begins with $(kz_0)^0$. From Eq. (15) we get $\nu_\mathrm{m}^{\mathrm{DP}} \sim T^5$ regardless of a type of surface; $\nu_\mathrm{m}^{\mathrm{PA}} \sim T^3$ for a semiconductor–vacuum system and $\nu_\mathrm{m}^{\mathrm{PA}} \sim T^5$ for a semiconductor covered with a thin metal film. Thus, for $T \ll T_0$, the power law for the temperature dependence of the near-surface relaxation rate $\nu_\mathrm{m}^{\mathrm{DP}}$ coincides with that of the QW in the bulk of semiconductor. At the same time, the temperature dependence of $\nu_\mathrm{m}^{\mathrm{PA}}$ for near-surface scattering agrees with the bulk value in the case of a semiconductor–vacuum system and differs (decreases faster with decreasing temperature) for a semiconductor covered with a thin metal film. Such a decrease in the PA scattering rate stems from a node of the piezoelectric interaction at semiconductor-metal interface due the boundary condition of Eq. (7).

In the foregoing consideration, a screening of the scattering potentials was not taken into account. The screening factor may be taken into account dividing the matrix elements $M_j(\omega, q)$ in Eq. (15) by a dielectric function $\varepsilon(w, q)$. For a 2DEG placed in a semi-bounded crystal contacting with vacuum [7]

$$\varepsilon(\omega, q) = 1 + \frac{2\pi e^2}{q\varepsilon_0} \Pi(\omega, q) F_\mathrm{s}, \tag{16}$$

where Π is the polarization function and

$$F_\mathrm{s} = \int_0^\infty \mathrm{d}z \psi^2(z) \int_0^\infty \mathrm{d}z' \psi^2(z') \times \left[\mathrm{e}^{-q|z-z'|} + \frac{\varepsilon_0 - 1}{\varepsilon_0 + 1} \mathrm{e}^{-q(z+z')} \right]. \tag{17}$$

It is worth mentioning that $\varepsilon(w, q)$ depends on a QW position z_0. For low temperature when $q \ll 2k_\mathrm{F}$, we have $\Pi = m^*/\pi\hbar^2$. If QW is placed near the surface so that $qz_0 \ll 1$, Eq. (16) reduces to

$$\varepsilon \simeq \frac{4\varepsilon_0}{(\varepsilon_0 + 1) a_\mathrm{B} q}, \tag{18}$$

where $a_\mathrm{B} = \varepsilon_0 \hbar^2 / m^* e^2$ is the effective Bohr radius. Since $\varepsilon_0 \gg 1$, the dielectric function of Eq. (18) which describes the near-surface screening appears to be two times larger than that for screening in the bulk semiconductor. Dividing the matrix elements $M_j(\omega, q)$ in Eq. (15) by ε from Eq. (18), in the limit of $T \to 0$ we get the following dependences: $\nu_\mathrm{m}^{\mathrm{DP}} \sim T^7$

and $\nu_m^{PA} \sim T^5$. These power laws agree with known results for a 2DEG in the bulk.

For a semiconductor in contact with metal, the screened potential obeys the boundary condition of Eq. (7). In this case, the factor $(\varepsilon_0 - 1)/(\varepsilon_0 + 1)$ in Eq. (17) is replaced with (-1). In the limit of small q, F_s becomes proportional to q, and $\varepsilon(\omega, q)$ of Eq. (16) appears to be independent of q. So, for a 2DEG placed near a metalized surface, screening does not alter the power laws of the momentum relaxation rates: ν_m^{DP} and ν_m^{PA} decrease as T^5. For finite T and z_0, the calculation of the temperature and distance dependences of the relaxation rates with screening included, will be presented separately.

In our calculations we have neglected the differences in the elastic, piezoelectric, and dielectric properties of the materials which form the heterostructure and have used only GaAs parameters. Calculations were carried out for the following parameters: $s_l = 5.2 \times 10^5$ cm/s, $s_t = 3.0 \times 10^5$ cm/s, $s_R = 2.77 \times 10^5$ cm/s, $\rho = 5.3\,\mathrm{g/cm^3}$, $\varepsilon_0 = 12.5$, $m^* = 0.067 m_0$, $D = 8$ eV, and $e_{14} = 0.16\,\mathrm{C/m^2}$. To extract the surface effect, we calculated the ratio $\nu_m(z_0)/\nu_m^b$, where the normalization function, ν_m^b, was determined as $\nu_m\,(z_0 \to \infty)$, and z_0 was the distance of the center of the QW from the surface. To derive ν_m^b from Eq. (15), one has to drop terms in the matrix element which are decaying exponentially and oscillating rapidly with functional forms of $\cos(2k_l z_0)$ and $\cos(2k_t z_0)$.

The temperature dependences of the normalized momentum relaxation rates of a 2DEG placed close to the surface of a semiconductor with vacuum and metal boundaries are shown in Figs. 3 and 4, respectively. The dependences on electron concentration, n_s, and QW width, d, are illustrated by curves 1–5. Fig. 3 shows that the proximity of the semiconductor–vacuum interface results in an enhancement of the electron–acoustic–phonon scattering. The effect is more pronounced for thin QWs with low electron densities. In the limit $T \to 0$, the ratio ν_m/ν_m^b reaches a maximum value which does not depend on n_s and d. This value is determined by PA scattering and depends only on the ratio of sound velocities, s_l/s_t.

As seen from Fig. 4, the temperature dependence of ν_m/ν_m^b in a semiconductor-metal system differs radically from that of a semiconductor bounding

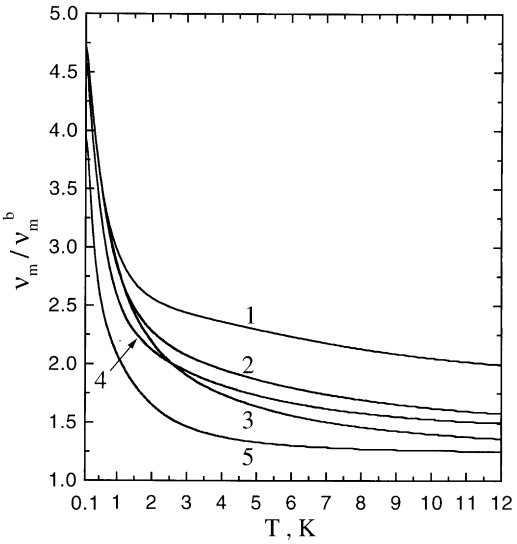

Fig. 3. Ratio of momentum relaxation rates to bulk values versus temperature for a semiconductor bounded by a vacuum. The distance from the crystal surface to the center of a QW equals the width of the well, d. Sheet electron concentrations n_s are (1) $10^{11}\,\mathrm{cm^{-2}}$; (2,4,5) $3 \times 10^{11}\,\mathrm{cm^{-2}}$; (3) $6 \times 10^{11}\,\mathrm{cm^{-2}}$; well widths, d: (1–3) 40 Å, (4) 50 Å, (5) 100 Å.

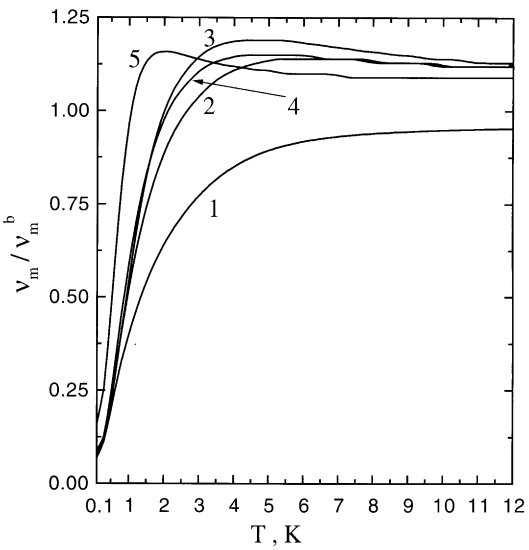

Fig. 4. Ratio of momentum relaxation rates to bulk values versus temperature for a semiconductor in contact with a metal. The notation used is the same as in Fig. 3.

a vacuum. The most significant result illustrated in Fig. 4 is a decrease in the near-surface scattering caused by the suppression of the PA coupling. In the limit as $T \to 0$, for both the PA and DP mechanisms, the partial contributions to ν_m depend on T as T^5. The bulk value of ν_m^b determined by the dominating PA scattering mechanism varies as T^3, and it follows that the ratio ν_m/ν_m^b scales as T^2 with decreasing temperature. It is worth noting that the characteristic temperature, T_c, which corresponds to the transition from DP-coupled to PA-coupled near-surface scattering differs from that of a QW in a bulk of semiconductor. For an electron sheet concentration of $n_s = 3 \times 10^{11}\,\mathrm{cm}^{-2}$ and for $d =$ 40 Å (curve 3), the critical temperature $T_c = 4.5\,\mathrm{K}$ for the bulk case. For the semiconductor–vacuum system, the same sheet density n_s, and $z_0 = 40$ Å, the calculated values of T_c differ substantially from the bulk results: $T_c = 11.1\,\mathrm{K}$. In the case of the semiconductor-metal system, the situation is more complicated. Only in the range of temperatures $T = 1.4$–$2\,\mathrm{K}$, does the PA-coupled near-surface scattering prevail over the DP-coupled scattering. For temperatures $T \gg 2\,\mathrm{K}$, deformation-potential scattering predominates. The dependences of the normalized momentum relaxation rates on the QW position have been calculated for the different electrical boundary conditions; the results are shown in Figs. 5 and 6. We see that the influence of a crystal surface has a long-distance character and the corresponding scale is larger for thin electron channels with low sheet concentrations.

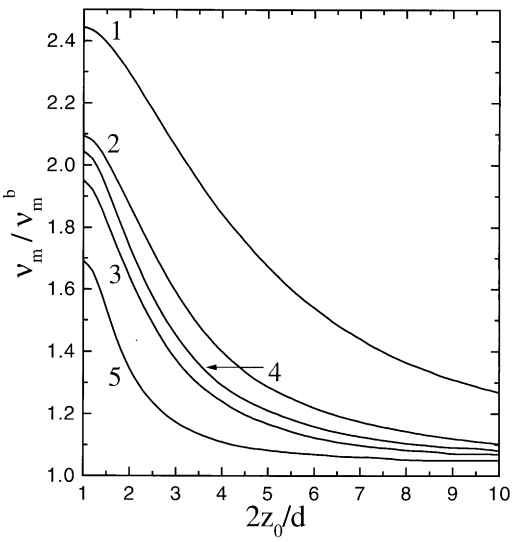

Fig. 5. Ratio of momentum relaxation rates to bulk values as a function of distance from the crystal surface to the center of QW, z_0, for a semiconductor bounded by a vacuum and with a temperature $T = 5$ K. The electron sheet concentrations, n_s, are: (1) $10^{11}\,\mathrm{cm}^{-2}$; (2,4,5) $3 \times 10^{11}\,\mathrm{cm}^{-2}$; (3) $6 \times 10^{11}\,\mathrm{cm}^{-2}$; well widths d: (1–3) 40 Å, (4) 50 Å, (5) 100 Å.

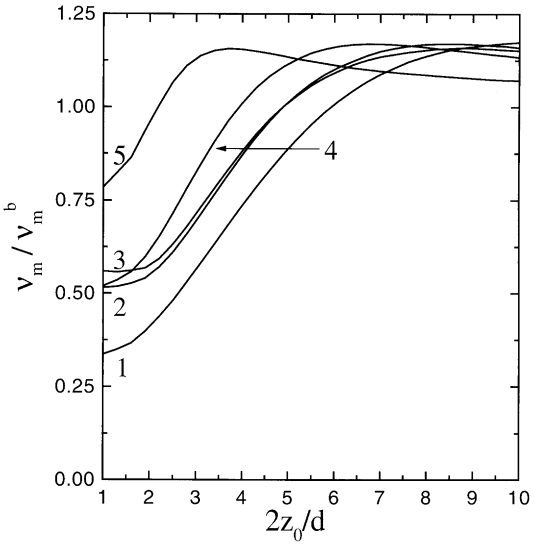

Fig. 6. Ratio of momentum relaxation rates to bulk values as a function of a distance from the crystal surface to the center of QW, z_0, for a semiconductor in contact with a metal film and a temperature $T = 1$ K. The notation used is the same as in Fig. 5.

4. Conclusions

We have studied the role of surface effects as they influence contribution of acoustic-phonon scattering to the momentum relaxation rate of a 2DEG. The scattering of the electrons via both the deformation and the piezoelectric interactions were taken into account. The peculiarities of the near-surface scattering originate from two sources: modification of the acoustic-phonon modes caused by the stress-free crystal surface and dependence of the phonon-induced piezoelectric potential on the dielectric properties of a medium in contact with the semiconductor. The first aspect of the problem has been solved using a complete set of the phonon

modes existing in a semi-bounded crystal. To study the boundary effect on the piezoelectric potential, the cases of a semiconductor in contact with vacuum and with a thin metal film have been considered. The results presented in this paper demonstrate a pronounced difference in the momentum relaxation rate of electrons in a narrow QW placed near a surface and the momentum relaxation rate of electron in a QW placed in the bulk of a semiconductor. It is shown that the relaxation rate is enhanced near the boundary with the vacuum; the relaxation rate is suppressed near a surface covered by a thin metal film.

Acknowledgements

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