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Hot-electron conduction in $Hg_{0.8}Cd_{0.2}Te$ in the extreme magnetic quantum limit at low temperatures

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A theoretical analysis of hot-electron transport in narrow-gap semiconductors under a large magnetic field is developed at low temperatures. The energy relaxation time and the electric-field dependence of the drift velocity of hot electrons in $Hg_{0.8}Cd_{0.2}Te$ at 1.5 K in the extreme quantum limit, i.e., when all the electrons occupy the lowest Landau subband, are calculated. The model includes scattering by the acoustic phonons and other complexities such as the band nonparabolicity, free-carrier screening, nonequipartition of phonons, and the level broadening due to electron-impurity interactions. The theoretical results are compared with the experimental data on the energy relaxation time in $Hg_{0.8}Cd_{0.2}Te$ for magnetic flux densities of 4 and 6 T. When the equilibrium phonon distribution is considered, the theoretical values are found to be several times less than the experimental values. The discrepancy is attributed to the nonequilibrium phonons or "hot phonons" that exist under hot-electron conditions. Including the hot-phonon effects, the theoretical results are in satisfactory agreement with the experimental data, giving a reasonable estimate of the phonon life time. The model has been applied to calculate the electric field dependence of the drift velocity of hot electrons in $Hg_{0.8}Cd_{0.2}Te$ at 1.5 K in the extreme quantum limit by assuming scattering by the nonequilibrium acoustic phonons and ionized impurities, and a satisfactory agreement with the experimental results is obtained.

I. INTRODUCTION

Hot-electron transport in semiconductors is a lively topic in view of the application in solid-state devices. Recently, Nimtz and Stadler[1] have reported experimental data on the energy relaxation time in $Hg_{0.8}Cd_{0.2}Te$ at 1.5 K in the extreme quantum limit (EQL). A theoretical analysis of these data is important for an insight into the basic energy loss mechanism controlling the hot-carrier kinetics in this technologically important semiconductor at low temperatures.

In this paper, a theoretical analysis of the energy relaxation of hot electrons at 1.5 K in the EQL is done. The model includes scattering by acoustic phonons together with the nonparabolicity of the conduction band, the screening of free carriers, and the broadening of the Landau level due to electron-impurity interactions. The equipartition approximation for the phonons is not made. The theoretical results are compared with the experimental results for the magnetic flux densities of 4 and 6 T.

When we take into consideration the thermal equilibrium phonons, the theoretical values of the energy relaxation time are found to be much less than the experimental values. We attribute this discrepancy to the nonequilibrium phonons or "hot phonons" that exist under hot-electron conditions. The emission of phonons by hot electrons and their reabsorption slow down the carrier cooling rate.[2] In order to explain the experimental data, the hot-phonon effect with reasonable values of the phonon lifetime for acoustic phonon scattering is considered. It is found that the incorporation of nonequilibrium acoustic phonons brings the theoretical results on the energy relaxation time in agreement with the experimental data giving reasonable values of phonon lifetime.

The success of nonequilibrium acoustic phonons in explaining the energy relaxation time has prompted us to investigate theoretically the electric-field dependence of drift velocity of hot electrons in $Hg_{0.8}Cd_{0.2}Te$ at 1.5 K in the EQL including scattering by nonequilibrium acoustic phonons and ionized impurities. The calculation shows a satisfactory agreement with recent experimental data of Stadler.[3]

II. DETAILS OF THE ANALYSIS

We assume a semiconductor with a nonparabolic, spherically symmetric conduction band subjected to a high magnetic flux density B along the z direction forcing the electrons to occupy the lowest spin-split Landau level, i.e., achieving the EQL condition. The energy dispersion relation for these electrons in the conduction band can be written as[4]

$$E = -\frac{E_g}{2} + \frac{E_g a_0}{2} + \frac{\hbar^2 k_z^2}{2m^* a_0}, \tag{1}$$

where E is the electron energy, k_z is the z component of the electron wave vector, $\hbar$ is the reduced Planck's constant, m^* is the band-edge effective mass, E_g is the band gap, and

$$a_0 = \left[1 + \frac{2\hbar\omega_c}{E_g}\left(1 - |g|\frac{m^*}{2m_0}\right)\right]^{1/2}. \tag{2}$$

In Eq. (2), g is the spin spectroscopic g factor, m_0 is the electron free mass, and $\omega_c = eB/m^*$, e being the electronic

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charge. We assume a nondegenerate electron gas obeying the Maxwell–Boltzmann distribution characterized by an electron temperature T_e subjected to a heating electric field applied parallel to B. The above assumption will be justified later. The average power loss per electron through emission and absorption of acoustic phonons is expressed as[5–7]

$$P_{ac} = \frac{(m^* a_0)^{1/2}\omega_0}{\pi(2\pi k_B T_e)^{1/2}\hbar} \exp\left(-\frac{m^* a_0 s^2}{2k_B T_e}\right) \int_0^\infty q_\perp \, dq_\perp \int_0^\infty \frac{dq_z}{q_z} C|f(q)|^2 \times \exp\left(-\frac{l^2 q_\perp^2}{2} - \frac{\hbar^2 q_z^2}{8m^* a_0 k_B T_e} - \frac{m^* a_0 s^2}{2k_B T_e}\cdot\frac{q_\perp^2}{q_z^2}\right)\left[(N_R+1)\exp(-\gamma_e) - N_R \exp(\gamma_e)\right], \tag{3}$$

where k_B is the Boltzman constant, $\hbar\omega_0$ is the phonon energy, s is the sound velocity, $l = (\hbar/eB)^{1/2}$ is the Landau radius, N_R is the phonon occupation number, $C|f(q)|^2$ is the electron-phonon coupling term, $q_\perp$ and q_z are the transverse and the logitudinal components of the phonon wave vector $\mathbf{q}$, and $\gamma_e = \hbar\omega_0/2k_B T_e$.

The acoustic phonon energy can be written as

$$\hbar\omega_0 = \hbar s q = \hbar s q_\perp (1 + q_z^2/q_\perp^2)^{1/2}. \tag{4}$$

Phonons mostly contributing to electron scattering are found to have a certain characteristic wavelength which, in the magnetic quantum limit, is taken to be of the order of the radius of the lowest Landau level. Thus we have $q_\perp \sim 1/l = (\hbar/eB)^{-1/2}$ and q_z can be taken for an average thermal electron in the z direction. Therefore, Eq. (4) allows the acoustic phonon energy to be approximated by[8] $\hbar s/l$. For deformation potential acoustic scattering, we have[5]

$$C|f(q)|^2 = C_{ac}\left[q^5/(q^2+q_s^2)^2\right], \tag{5}$$

where

$$C_{ac} = E_1^2 \hbar/2\rho s, \tag{6}$$

E_1 it the acoustic deformation potential constant, ρ is the mass density, and q_s is the inverse screening length.

For piezoelectric coupling to acoustic phonons, we have

$$C|f(q)|^2 = C_{pz}\left[q^3/(q^2+q_s^2)\right], \tag{7}$$

where

$$C_{pz} = \hbar^2 e^2 e_{14}^2/2\rho u_p \epsilon_s^2, \tag{8}$$

e_{14} is piezoelectric modulus, u_p is the piezoelectric velocity, and ϵ_s is the permittivity of the material.

We first treat the case of thermal phonons for which the occupation number is given by Bose–Einstein statistics. We use Eqs. (5)–(8) in Eq. (3) and make the following substitutions:

$$\gamma = \frac{\hbar^2 s^2}{8(k_B T_e l)^2}, \quad \tfrac{1}{2}(lq_\perp)^2 = u,$$
$$\beta = 8m^* a_0 l^2 k_B T_e/\hbar^2,$$
$$\alpha_s = q_s^2 l^2, \quad \text{and} \quad \upsilon = \hbar^2 q_z^2/(8m^* a_0 k_B T_e).$$

For deformation potential scattering the power loss per electron is then given by

$$P_{ac} = W_{ac}\int_0^\infty f_{ac}(\upsilon)\exp(-\upsilon)\frac{d\upsilon}{\upsilon}, \tag{9}$$

where

$$W_{ac} = \frac{(m^* a_0)^{1/2} E_1^2 \omega_0 N_0}{8\pi(k_B T_e)^{1/2} l^3 \rho s}\left[\exp\left(\frac{\hbar\omega_0}{k_B T_L}\right) - \exp\left(\frac{\hbar\omega_0}{k_B T_e}\right)\right] \times \exp\left(-\frac{\hbar\omega_0 + m^* a_0 s^2}{2k_B T_e}\right) \tag{10}$$

and

$$f_{ac}(\upsilon) = \left(1+\frac{\beta\upsilon}{u_0}\right)^{5/2}\left(1+\frac{\gamma}{\upsilon}\right)^{-3/2}\left(1+\frac{\beta\upsilon+\alpha_s}{u_0}\right)^{-2}. \tag{11}$$

For piezoelectric scattering the power loss per electron is expressed as[5]

$$P_{pz} = W_{pz}\int_0^\infty f_{pz}(\upsilon)\exp(-\upsilon)\frac{d\upsilon}{\upsilon}, \tag{12}$$

where

$$W_{pz} = \frac{(m^* a_0)^{1/2} e^2 e_{14}^2 N_0\omega_0}{8\pi\rho u_p \epsilon_s^2 (k_B T_e)^{1/2} l}\left[\exp\left(\frac{\hbar\omega_0}{k_B T_L}\right) - \exp\left(\frac{\hbar\omega_0}{k_B T_e}\right)\right]\exp\left(-\frac{\hbar\omega_0 + m^* a_0 u_p^2}{2k_B T_e}\right) \tag{13}$$

and

$$f_{pz}(\upsilon) = \left(1+\beta\frac{\upsilon}{u_0}\right)\left(1+\frac{\gamma}{\upsilon}\right)^{-1/2}\left(1+\frac{\beta\upsilon+\alpha_s}{u_0}\right)^{-2}. \tag{14}$$

Here the integration over u is simplified by replacing u in the slowly varying function

$$\left(1+\frac{\beta\upsilon}{2u}\right)^{5/2}\left(1+\frac{\beta\upsilon+\alpha_s}{2u}\right)^{-2}$$

by $u_0/2$. Since $u_0 = l^2 q_\perp^2 = 1$, this function, to a good approximation,[7] can be taken outside the integral over u.

The quantities P_{ac} and P_{pz}, given by Eqs. (9) and (12), show a divergence at $\upsilon = 0$ or $q_z = 0$. We get around this problem by incorporating the broadening of the Landau level due to electron-impurity interactions. This corresponds to taking the lower limit of the integrals in Eqs. (9) and (12) as $E_c/4k_B T_e$, where E_c is the cutoff energy[9] varying with the magnetic flux density B as $B^{2/3}$. Also, we use the approximations that $f_{ac}(\upsilon)$ and $f_{pz}(\upsilon)$ are slowly varying functions of υ so that they can be taken outside the integrals with υ replaced by its optimum value of unity. This corresponds to taking $\hbar^2 q_z^2/2m^* a_0 = 4k_B T_e$, which is approximately the upper limit of q_z. We can then write

$$P_{ac} = W_{ac} f_{ac} \ln[(4k_B T_e/E_c)\exp(-c_e)] \quad (15)$$

and

$$P_{pz} = W_{pz} f_{pz} \ln[(4k_B T_e/E_c)\exp(-c_e)] \quad (16)$$

where c_e is Euler's constant and the functions $f_{ac}(v)$ and $f_{pz}(v)$ are evaluated at $v = 1$.

The energy relaxation time τ_e is determined from the relationship

$$P_{ac} + P_{pz} = \tfrac{1}{2}k_B(T_e - T_L)/\tau_e . \quad (17)$$

The factor 1/2 on the right-hand side of Eq. (11) accounts for one translational degree of freedom due to Landau quantization.

Numerical calculations are done with the values of the material constants[10] of $Hg_{0.8}Cd_{0.2}Te$ given in Table I. The cutoff energy E_c related to level broadening is not accurately known. However, it must be comparable to the thermal energy and is taken to be 0.1 meV at $B = 4$ T. The calculated results are not sensitive to the uncertainty in the value of E_c since E_c enters through a logarithmic term.

The calculated values of τ_e are presented in Fig. 1 along with the experimental values. Electron temperatures between 2.5 and 3.5 K are considered as the specific heat data[1] show that the carrier system behaves like a Maxwell–Boltzmann gas over this range of temperatures. The application of the magnetic field with the band nonparabolicity makes the degeneracy effects insignificantly small as the Fermi level is about $2k_B T_e$ below the band edge for a carrier concentration of about 10^{14} cm^{-3} found in the experimental sample.

The quantitative agreement between the experimental and the theoretical values of τ_e is found to be poor. The experimental values of τ_e are much higher than the theoretical values. We thus find that we cannot explain the experimental results by including complexities such as band nonparabolicity, nonequipartition of phonons, screening of free carriers, and Landau level broadening by considering thermal phonons alone.

We also find from Fig. 1 that the deformation potential coupling to acoustic phonons is primarily responsible for the energy loss of hot electrons. In the following analysis, including nonequilibrium phonons, we therefore consider only the deformation potential coupling and neglect the piezoelectric coupling for simplicity.

TABLE I. Material parameters of $Hg_{0.8}Cd_{0.2}Te$ used for the calculation given by Dornhaus and Nimtz (Ref. 10).

Effective mass (m^*)	$0.0065m_0$
Band gap (E_g)	70 meV
Mass density (ρ)	7.654×10^3 kg m^{-3}
Lande's g factor $\lvert g\rvert$	90 ($B = 4$ T)
	80 ($B = 6$ T)
Acoustic velocity (s)	3.017×10^3 ms^{-1}
Acoustic deformation potential (E_1)	9.0 eV
Piezoelectric modulus (e_{14})	0.0335 C m^{-2}
Piezoelectric velocity (u_p)	1.948×10^3 ms^{-1}
Dielectric constant (K_0)	18

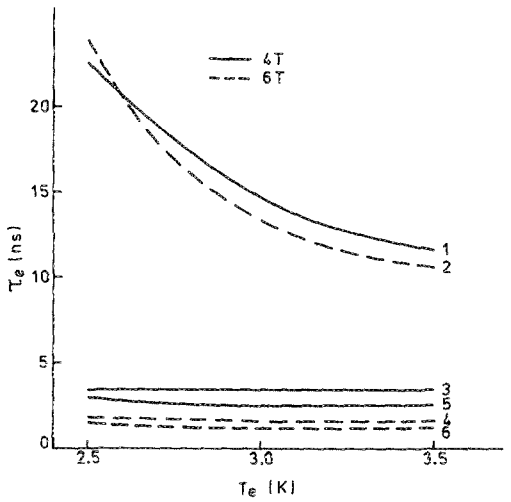

FIG. 1. Variation of τ_e with T_e at $T_L = 1.5$ K. Curves 1 and 2 show the experimental values for $B = 4$ and 6 T, respectively. Curves 3 and 4 show the calculated values for deformation potential coupling assuming thermal phonons for $B = 4$ and 6 T, respectively. Curves 5 and 6 show the same for the additional effect of piezoelectric coupling assuming thermal phonons for $B = 4$ and 6 T, respectively.

III. NONEQUILIBRIUM PHONONS

The rate of generation of phonons due to hot electrons in the steady state can be expressed as[2]

$$\left(\frac{\partial N_R}{\partial t}\right)_e = \frac{N_R - N_0}{\tau_p}, \quad (18)$$

where N_R and N_0 are, respectively, the nonequilibrium and thermal phonon occupation numbers. The right-hand side of Eq. (18) gives the rate of loss of phonons with a lifetime τ_p. The transfer of power per electron is given by

$$P_{ac} = -\frac{1}{nV}\sum_q \hbar\omega_0\left(\frac{\partial N_R}{\partial t}\right)_e = \frac{\hbar\omega_0}{4\pi^2 n}\left(\frac{\partial N_R}{\partial t}\right)_e \int_0^{q_{\perp m}} q_\perp\, dq_\perp \int_0^{q_{zm}} dq_z . \quad (19)$$

Here n is the density of carriers, V is the crystal volume, $q_{\perp m}^2 = l^{-2}$, and $\hbar^2 q_{zm}^2/2m^*a_0 = 4k_B T_e$. Following our previous approximations we shall replace ω_0 by s/l in Eq. (19). Hence

$$\left(\frac{\partial N_R}{\partial t}\right)_e = \frac{4\sqrt{2}\pi^2 n l^2}{\omega_0(m^* a_0 k_B T_e)^{1/2}} P_{ac} . \quad (20)$$

Using Eqs. (3), (18), and (20) we are led to

$$N_R = \frac{a\alpha\exp(-\gamma_e) + N_0/\tau_p}{a\alpha[\exp(\gamma_e) - \exp(-\gamma_e)] + \tau_p^{-1}}, \quad (21)$$

where

$$a = \frac{4\sqrt{2}\pi^2 n l^2}{\omega_0(m^* a_0 k_B T_e)^{1/2}},$$

$$\alpha = W_{ac1} f_{ac} \ln[\exp(-c_e)(4k_B T_e/E_c)],$$

and

$$W_{ac1} = \frac{(m^* a_0)^{1/2} E_1^2 \omega_0}{8\pi(k_B T_e)^{1/2} l^3 \rho s}\exp\left(-\frac{m^* a_0 s^2}{2k_B T_e}\right).$$

Under hot-phonon conditions Eq. (11) can be written as

$$\alpha[(N_R+1)\exp(-\gamma_e)-N_R\exp(\gamma_e)] = \frac{(1/2)k_B(T_e-T_L)}{\tau_e}. \quad (22)$$

To calculate τ_e, the values of τ_p must be known. But exact values of τ_p are not known accurately because of the complex nature of the phonon relaxation mechanism.[11,12] We therefore think it judicious to examine whether the experimental values of τ_e can be fitted with reasonable values of τ_p. The calculated values of τ_p required for such a fit are shown in Table II.

We find that the required values of τ_p lie in the range 32–112 ns. Such values are not unreasonable according to the discussions by Aronzon and Meilikhov[12] and Askerov.[13] The values of τ_p of this order have also been considered by Jay-Gerin[11] in evaluating the phonon drag contribution to the thermoelectric power in EQL in GaAs at low temperatures.

As a rough approximation, the contribution of the boundary scattering to τ_p can be taken as $(L/2s)$, where L is the shortest transverse dimension of the sample. For the given sample, τ_p, calculated in this manner, turns out to be 100 ns, which falls within the range of the values of τ_p found in our analysis. However, the mismatch factor with boundaries may substantially increase the value of τ_p estimated from the boundary scattering.[14] The values of τ_p found in our analysis would then be less than that determined by the boundary scattering including the mismatch factor. Other processes, such as phonon-phonon interaction, therefore seem to control the decay of phonons in the present situation.

The results of Table II show that for a given magnetic field, τ_p decreases with increasing T_e. With increased heating of the electron gas, phonon-phonon interaction seems to become stronger, accounting for such behavior. The hot-phonon effect is thus believed to control the cooling of the carriers in $Hg_{0.8}Cd_{0.2}Te$. It will be interesting to investigate the electric-field dependence of the drift velocity in this material at 1.5 K in EQL, including scattering by nonequilibrium acoustic phonons and ionized impurities.

IV. THE DRIFT VELOCITY

In order to calculate the drift velocity for hot electrons, we must know the momentum relaxation rate of the carriers in addition to the energy loss by nonequilibrium acoustic phonons. The relaxation rate of the carriers is

TABLE II. Calculated results (B = magnetic flux density, T_e = electron temperature, τ_p = phonon life time).

$B(T)$	T_e(K)	τ_p(ns)
4	2.5	62.4
4	3.0	41.0
4	3.5	32.1
6	2.5	111.5
6	3.0	66.3
6	3.5	55

$$\frac{1}{\tau}=\frac{1}{\tau_a}+\frac{1}{\tau_i}, \quad (23)$$

where τ_a and τ_i are the relaxation times for the acoustic and the ionized impurity scattering, respectively. The expressions for τ_a and τ_i are derived here following the line of arguments given earlier. The final expressions read

$$\frac{1}{\tau_a}=\frac{(m^*a_0)^{1/2}}{4\pi^{1/2}\hbar^2 l^3}C_{ac}\left(1+\frac{\beta v}{u_0}\right)^{5/2}\left(1+\frac{\beta v+\alpha_s}{u_0}\right)^{-2} \times\left(\frac{N_R+1}{\sqrt{E-\hbar\omega_0}}+\frac{N_R}{\sqrt{E+\hbar\omega_0}}\right), \quad (24)$$

and

$$\frac{1}{\tau_i}=\frac{N_i(el)^4}{16\pi\hbar l^2\epsilon_s^2}\left(\frac{2m^*a_0}{\hbar^2}\right)^{1/2}E^{-1/2} \times\int_0^\infty\left(u+\frac{\beta v+\alpha_s}{2}\right)^{-2}e^{-u}\,du. \quad (25)$$

Here N_i represents the concentration of ionized impurities, considered to be equal to the concentration of free electrons.

Equations (23), (24), and (25), along with the distribution function, allow a determination of the electron mobility, which, coupled with Eq. (17), gives the drift velocity and the electric field. Numerical calculations are performed using the parameter values of $Hg_{0.8}Cd_{0.2}Te$ given in Table I. The values of τ_p that fit the theoretical values of the energy relaxation time to the experiments[3] are used. In order to compare with experiments, the field F is normalized by its value F_n corresponding to the electron temperature of 2.5 K at which the system begins to behave like a nondegenerate gas.[1,3] The values of F_n for $B=4$ and 6 T are 8 and 12 V m^{-1}, respectively. We normalize the drift velocity V_d by the value V_{dn} which it attains at the applied flux density of 4 T at $T_e=2.5$ K. The value of V_{dn} is 4.8 ms^{-1}. The plot of V_d/V_{dn} vs F/F_n for $B=4$ and 6 T is given in Fig. 2. The use

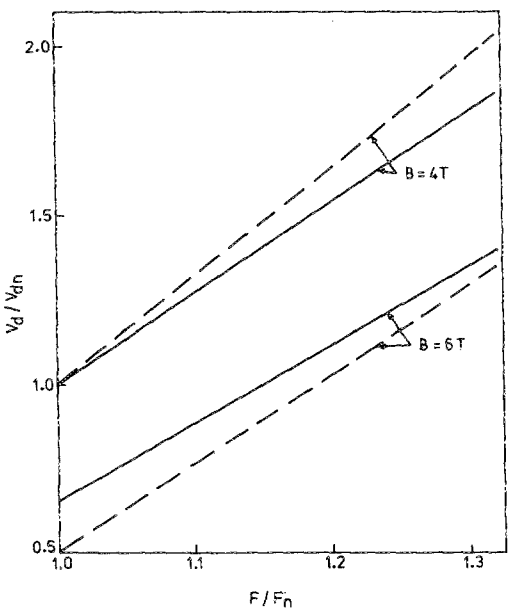

FIG. 2. Variation of the normalized drift velocity (V_d/V_{dn}) with the normalized elecric field (F/F_n). The solid curves give the theoretical values while the dashed curves give the experimental data of Stadler (Ref. 3).

of the normalized variables permits a comparison of the theoretical values of V_d/V_{dn} obtained here with the experimental values of the corresponding current ratios. The possible uncertainty in the values of the carrier concentration will not affect the comparison due to the normalization.

The carrier cooling rate is slowed down by the nonequilibrium phonons (see Sec. III), so that the electron temperature rises rapidly with increasing field. The effect of ionized impurity scattering which is basically Coulombic in nature is reduced by the carrier heating. The enhancement of the mobility due to the weakening of the ionized impurity scattering with increasing field is, however, offset by the acoustic phonon scattering, the effect of which increases with carrier heating. Therefore, the drift velocity increases almost linearly with electric field, as shown in Fig. 2. The variation of the drift velocity with the field is found to agree with experiments[3] for both $B = 4$ and 6 T. The maximum deviation of the theoretical values from the experimental ones is about 20% which can be accounted for by uncertainties in the parameter values.

In conclusion, our model of hot-electron transport in $Hg_{0.8}Cd_{0.2}Te$ under extreme magnetic quantum limit at low temperatures can explain not only the experimental values of the energy relaxation time but also the experimentally observed variation of drift velocity with electric field.

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