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Electron–phonon coupling and phonon-drag thermopower of a very low mobility 2DEG

R. Fletcher^{a, *}, M. Tsaousidou^b, P.T. Coleridge^c, Y. Feng^c, Z.R. Wasilewski^c^aPhysics Department, Queen's University, Kingston, Canada K7L3N6^bSection of Solid State Physics, University of Athens, Zografos 157 84, Greece^cMicrostructural Sciences, National Research Council, Ottawa, Canada K1A0R6

Abstract

It has been predicted that the phonon drag thermopower will be enhanced in very low mobility samples. The effect occurs when $ql < 1$ where q is the phonon wave number and l the electron mean free path. We present experimental data and detailed numerical calculations which confirm this prediction. © 2002 Elsevier Science B.V. All rights reserved.

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Under normal circumstances the electron–phonon (e–p) coupling in 2D systems is both predicted and found to be independent of sample mobility μ . However, when $ql < 1$, where q is the phonon wave number and l the electron mean free path, the coupling is predicted to increase [1,2]. This situation is reached only with low mobility samples at low temperatures. The reason for the change in behaviour in this limit is that the diffusivity of the electrons becomes so low that the screening of the potential associated with the acoustic wave becomes less effective. The effect has previously been investigated using the *energy loss rate* of electrons [1,3]. The present experiments use phonon drag thermopower, S^g , to probe the *momentum relaxation rate* of the electrons due to e–p scattering.

The measurements were made on a gated, GaAs/GaAlAs single heterojunction. To accentuate

electron-impurity scattering, there was no spacer between the Si-doped layer and the channel, and the latter was also intentionally doped with $1 \times 10^{22} \text{ m}^{-3}$ Be. The thermopower, S , and conductivity, σ , were measured at various values of electron density, n , in the temperature range of $0.3 \text{ K} < T < 4.2 \text{ K}$. The sample characteristics are shown in Table 1. We have used $q = k_B T / \hbar v_t$ and the mobility at 4.2 K to estimate ql/T , where v_t is the transverse sound velocity $\sim 3000 \text{ m/s}$.

At all values of n , σ exhibited a $\ln T$ dependence indicating weak localization (WL) effects. At low n this was very pronounced such that by $n = 1.79 \times 10^{15} \text{ m}^{-2}$, the sample resistance increased by about a factor of 3 between 4.2 and 0.3 K.

Both DC and AC methods were used for S and gave consistent results. The combination of DC and a high resistance sample gives a spurious offset voltage due to the input bias current of the Keithley 182 voltmeter used as a detector. This offset was first minimized by using a Hall bar with a small value of the geometric ratio l/w to keep the sample resistance low. It was

* Corresponding author. Tel.: +1-613-533-2705; fax: +1-613-533-6463.

E-mail address: fletcher@physics.queensu.ca (R. Fletcher).

Table 1
Sample characteristics

n (10^{15} m^{-2})	μ (m^2/Vs) 4.2 K	ql/T (K^{-1})
1.79	0.17	0.5
2.00	0.25	0.8
2.45	0.43	1.5
3.09	0.76	3
4.18	1.4	6
5.26	2.0	10

then eliminated from the results by taking data with the potential probes on the sample connected in the two possible ways and averaging. With AC there was no such offset but the AC amplitudes (which were weakly frequency dependent in the range 1–8 Hz but independent of n) were normalised using the DC data.

At low T (in the Bloch limit) when piezoelectric scattering is the dominant e–p scattering mechanism, theory predicts that $S^g \propto AT^4/n^{3/2}$ [2,4] for clean samples ($ql \gg 1$), where the prefactor is known and A is the phonon mean free path. Experiments, both present and previous, are in good accord with this relation. For dirty samples ($ql \ll 1$) it is predicted that $S^g \propto AT^3/D$ [2] at low T where D is the electron diffusion coefficient. A depends only on the properties of the substrate, not the 2DEG, and we have eliminated this dependence of S^g on A by working with the ratio S/λ where S is the measured thermopower and λ is the thermal conductivity of the substrate given by $\lambda \propto AT^3$, the constant again being known. The quantity S^g/λ has the dual advantages of being independent of the details of the particular substrate and so can be readily compared for different samples, and can also be measured more accurately because it does not involve ∇T .

The experimental situation is somewhat complicated by the presence of a diffusion contribution to the S of the form $S^d \propto T$. Thus, at low T in the two limits we expect S/λ to behave like $AT^\gamma + CT^{-2}$ where A and C are constants and $\gamma = 1$ for clean samples and 0 for dirty samples (assuming D is independent of T). The experimental data do indeed have this form and so S^g and S^d were readily separated. The sample with $n = 2.45 \times 10^{15} \text{ m}^{-2}$ is not clearly in either limit so that S^g does not have a simple functional form. In this case, we simply subtracted a term CT^{-2} from S/λ to find S^g/λ , where the value of C was interpo-

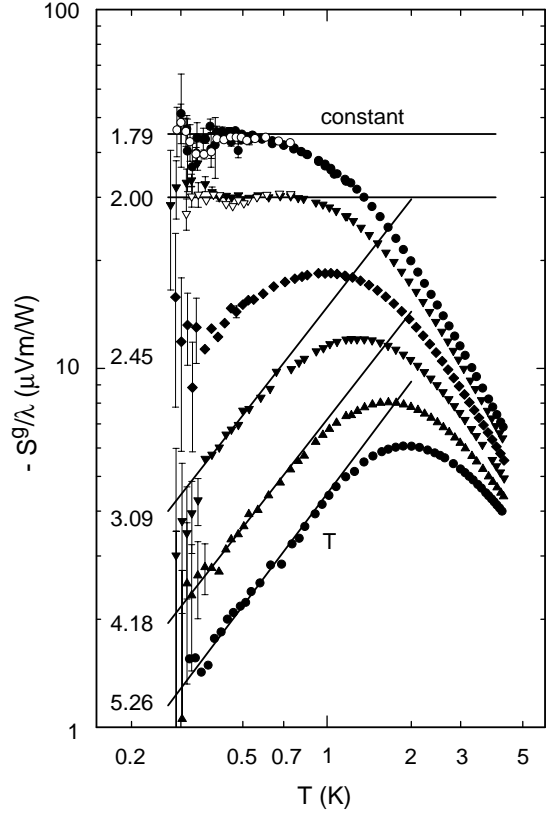


Fig. 1. Experimental data on S^g/λ as a function of T for various densities (given in units of 10^{15} m^{-2}). The solid and open symbols are DC and AC measurements, respectively. The lines through the data at higher densities are $\propto T$.

lated from the others. Note that we see no evidence that S^d diverges at a finite value of n as was recently found with Si-MOSFETs [5]. In fact, S^d/T has a rather weak dependence on n and for the two dirty samples S^d is only a small fraction of S even at our lowest temperature. After removing the diffusion contribution from all the data we find the results for S^g/λ shown in Fig. 1. Just as expected, we see that for the three clean samples $S^g/\lambda \propto T$ at low temperatures and for the two dirty samples S^g/λ tends to a constant.

We now turn to a detailed calculation to compare with the above experimental results. The calculation is based on the standard Cantrell–Butcher (CB) formalism [6]. Here we use the refined CB formula which includes both non-degeneracy and screening effects [7]. For the purposes of our analysis the expression of S^g

is written in the following form:

$$S^g = -\frac{\Lambda \hbar^2}{16\pi^2 n |e| k_B T^2 \rho} \sum_{i=1}^3 \int_0^\infty \int_0^\infty dq dq_z \times \frac{\Xi_{\text{eff}}^2}{|\varepsilon(\omega, q)|^2} \frac{Z(q_z) Q q^3}{\sinh^2(\hbar\omega/2k_B T)} \text{Im} \chi(q, \omega), \quad (1)$$

where ρ is the density of the material, ω is the frequency of phonons with wave vector $\mathbf{Q} = (\mathbf{q}, q_z)$ and polarization i (\mathbf{q} is the component in the plane of the 2DEG), $Z(q_z)$ is a form factor that accounts for the finite thickness of the 2DEG, and Ξ_{eff} is the ‘effective’ acoustic potential. In GaAs, Ξ_{eff} accounts for both deformation potential and piezoelectric coupling. Explicit expressions for Ξ_{eff} for the longitudinal and transverse branches are given by Lyo [8]. $\varepsilon(\omega, q)$ is the dynamic 2D dielectric function for which we have used the expression given in Refs. [2,9]. Finally, $\text{Im} \chi(q, \omega)$ is the imaginary part of the density polarization function

$$\text{Im} \chi(q, \omega) = \frac{2\pi}{A} \sum_{\mathbf{k}} [f^0(\varepsilon_{\mathbf{k}}) - f^0(\varepsilon_{\mathbf{k}+\mathbf{q}})] \times \delta(\hbar\omega + \varepsilon_{\mathbf{k}} - \varepsilon_{\mathbf{k}+\mathbf{q}}), \quad (2)$$

where A is the area of the 2DEG, \mathbf{k} is the 2D electron wave vector, $\varepsilon_{\mathbf{k}} = \hbar^2 k^2 / 2m^*$ is the electronic energy and $f^0(\varepsilon_{\mathbf{k}})$ is the Fermi–Dirac function. In the clean regime at finite temperatures we evaluated $\text{Im} \chi(q, \omega)$ numerically. $\text{Im} \chi(q, \omega)$ can be written in an analytic form, $\text{Im} \chi(q, \omega) = (m^* / \pi \hbar^2)(\omega / v_F q)$, only when $T \rightarrow 0$ [10]. In the dirty limit the density polarization function is characterized by the diffusion coefficient $D = v_F l / 2$ [11] and $\text{Im} \chi(\omega, q) = (m^* / \pi \hbar^2)(D q^2 / \omega) / [(D q^2 / \omega)^2 + 1]$ [2,11]. This difference has very pronounced effects on S^g .

We have calculated S^g by using Eq. (1) and the standard material parameters for GaAs [8]. The values of the sound velocities that we used are $v_l = 4969$ m/s and $v_t = 2755$ m/s for the longitudinal and the transverse modes, respectively. These were obtained by averaging $1/v_i^5$ over the [1 0 0], [1 1 0], and [1 1 1] crystallographic directions [4]. The thermal conductivity was calculated by using the expression $\lambda = (2\pi^2 k_B^4 / 45 \hbar^3) \Lambda T^3 \sum_i 1/v_i^2$. Similar averaging of $1/v_i^2$ over the three high symmetry directions gives $\lambda = 3890 \Lambda T^3$ W/mK [4].

The calculated results for the ratio S^g/λ are shown as solid lines in Fig. 2. For these curves we have

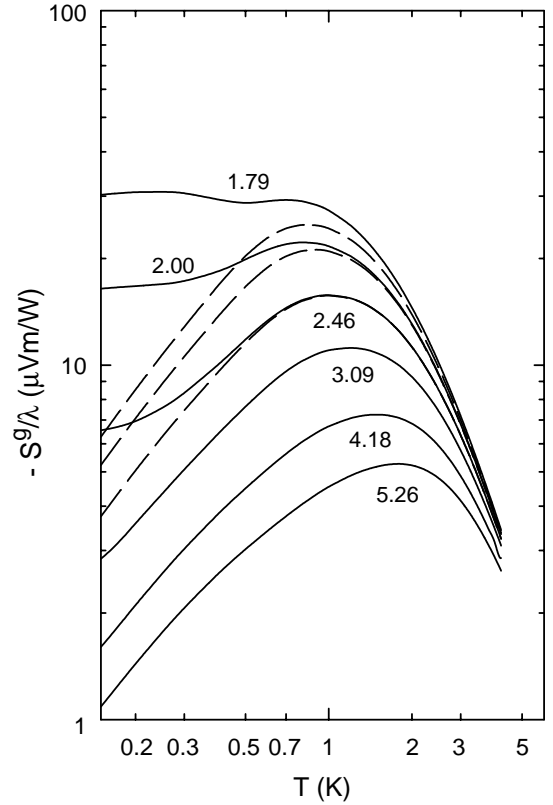


Fig. 2. The solid lines were calculated using the full theory for the same data as shown in Fig. 1. The dashed lines for the three lowest mobility samples were calculated assuming the clean limit to show the differences.

used the electron mean free path l as obtained from the Boltzmann part, σ_0 , of the conductivity by using the Einstein relation $\sigma_0 = (m^* e^2 / \pi \hbar^2) D$. To obtain σ_0 we needed to remove WL effects from the measured conductivity. Since WL effects are suppressed at high T and in the presence of a magnetic field, we have approximated σ_0 by the measured σ at $T = 4.2$ K and $B = 1$ T. Comparing Figs. 1 and 2 we see that the theoretical values of S^g/λ are in reasonably good agreement with the experimental values. We should mention that when l is estimated from the measured σ the calculated values of S^g/λ are far too large and have the wrong T -dependence at low T (not shown in the figure).

To get a clearer picture of how strongly S^g is enhanced in the dirty limit, Fig. 2 also shows calculated

values of S^g/λ (dashed lines) for the samples with $n = 1.79, 2.0$, and $2.46 \times 10^{15} \text{ m}^{-2}$ when the $ql < 1$ effect is ignored. We see that the absolute magnitude of S^g is considerably larger at low T in the dirty regime. This can be understood by examining the behaviour of the function $\text{Im } \chi(\omega, q)/|\varepsilon(\omega, q)|^2$ that appears in the RHS of Eq. (1) in the two regimes. For reasons of simplicity, we confine our interest to low temperatures where we can easily show that

$$\frac{\text{Im } \chi(\omega, q)}{|\varepsilon(\omega, q)|^2} \approx \begin{cases} (\pi\hbar^2/m^*)(\omega/v_F q)(1/V_{e-e}^2) & \text{clean limit,} \\ (\pi\hbar^2/m^*)(\omega/Dq^2)(1/V_{e-e}^2) & \text{dirty limit,} \end{cases} \quad (3)$$

where V_{e-e} is the 2D Fourier transform of the Coulomb interaction. Inspection of Eq. (3) shows that the function $\text{Im } \chi(\omega, q)/|\varepsilon(\omega, q)|^2$ is enhanced by the factor $2/ql$ in the dirty limit. This factor also explains why the power law of S^g is T^3 instead of the standard T^4 behaviour for clean samples.

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