Oscillatory Magnetothermopower and Resonant Phonon Drag in a High-Mobility 2D Electron Gas

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Experimental and theoretical evidence is presented for new low-magnetic-field ($B < 5$ kG) $1/B$ oscillations in the thermoelectric power of a high-mobility GaAs/AlGaAs two-dimensional (2D) electron gas. The oscillations result from inter-Landau-level resonances of acoustic phonons carrying a momentum equal to twice the Fermi wave number at $B = 0$. Numerical calculations show that both 3D and 2D phonons can contribute to this effect.

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Thermoelectric power (TEP) of a two-dimensional electron gas (2DEG) reflects the electron density of states, scattering dynamics, and electron-phonon interactions. Phonon scattering and consequently the TEP are particularly important in $GaAs/Al_xGa_{1-x}As$ heterostructures owing to the strong electron-lattice coupling in these materials. The TEP experiments in the following two limits of the magnetic field (B) have been widely pursued [1–8] and the phenomena in these regimes are relatively well understood [9–13]. At $B = 0$, the TEP (S_0) shows a power law, e.g., T^{3-4}, temperature dependence above $T \sim 0.3$ K, indicating that the phonon-drag mechanism dominates and the electron-diffusion effect is relatively weak. In a higher field, Shubnikov–de Haas (SdH) oscillation and quantum Hall effect have been observed in TEP [8]. The peak amplitudes of the TEP are found to be greater than the predicted values of the diffusion TEP by 2 orders of magnitude, but consistent with the phonon-drag TEP. Most of the above experimental data at $B = 0$ and at relatively high fields have been successfully explained quantitatively by theories based on the phonon-drag TEP.

On the contrary, little is known experimentally about the TEP in a weak magnetic field, where many Landau levels (LLs) are occupied by electrons with large quantum numbers $n \gg 1$ at the Fermi level, and electronic transport is generally treated semiclassically. This regime is characteristically distinct for the following reasons. (i) The Fermi wavelength is much shorter than the magnetic length $l_B = \sqrt{\hbar/eB}$ (i.e., $2k_F\ell_B \gg 1$), and a momentum selection rule governs the scattering of electron guiding centers. In particular, $2k_F$ scattering is strongly enhanced, giving rise to the low B resonance phenomena [14,15] in this regime. (ii) Acoustic phonons in GaAs have suitable energies, and in combination with (i) can effectively participate in inter-LL scattering. This is dramatically different from the regime of higher B, where intra-LL scattering dominates and inter-LL scattering is negligible at low temperatures [12]. Elastic intra-LL scattering is important for LL broadening.

In this Letter, we report low-B TEP oscillations observed in a high-mobility 2DEG in $GaAs/Al_xGa_{1-x}As$ heterostructures. The oscillations are periodic in $1/B$ with the B position of peaks proportional to $\sqrt{n_e}$, or to the Fermi wave number $k_F = \sqrt{2\pi n_e}$, where n_e is the electron density. Characteristically, such oscillations appear in the temperature range of 0.5–1 K, and their amplitude increases with T. It will be shown that the TEP oscillations result from inter-LL cyclotron resonance promoted by resonant absorption and emission of phonons carrying a $2k_F$ in-plane momentum and an energy $\hbar 2k_F u = \ell\hbar\omega_c$ equal to the integer multiple ($\ell = 1, 2, \ldots,$) of the cyclotron energy $\hbar\omega_c = \hbar eB/m^*$, where u is the phonon velocity, and m^* is the electron effective mass. 2D phonon modes propagating along the $GaAs/Al_xGa_{1-x}As$ interface and 3D phonons are responsible for the TEP oscillations. However, the contribution from 2D phonons is difficult to assess quantitatively because of the lack of electron-2D phonon interaction parameters. On the other hand, numerical calculations show that 3D phonons yield a substantial contribution to the oscillation and behave like 2D phonons for the $2k_F$ oscillation because q_z is restricted to small $q_z \ll 2k_F$ at low temperatures due to the phonon occupation factor.

The TEP experiments were performed using a sorption pumped ^{3}He cryostat and a superconducting solenoid, with the B axis always oriented perpendicular to the 2DEG plane. The samples used here are made from $GaAs/Al_{0.3}Ga_{0.7}As$ heterostructures grown by molecular beam epitaxy on the (001) GaAs substrate. At low temperature ($T \sim 1$ K), the n_e and the mobility, μ, can be varied using a red light-emitting diode (LED). Without LED (saturate LED) the n_e (in units of 10^{11} cm^{-2} throughout this paper) ~1.33 (2.03) and $\mu \sim 2 \times 10^6$ (3×10^6) cm^2/V s. To make TEP measurement, a specimen of dimension 8 mm $\times$ 2 mm was first cleaved from the wafer and a Hall bar mesa (width 0.5 mm) was then chemically wet etched from it by optical lithography. Heat sinking to a copper post (0.3 K) was achieved by indium soldering at one end of the specimen. A strain gauge was glued to the other end of the specimen by Ag paint and was used as a heater to create a temperature gradient (∇T) along the Hall bar direction. Two calibrated RuO_2 chip sensors were glued by Stycast 2850FT [16] on the back of the specimen and used to measure the ∇T along the sample. Electrical leads were made with 38 gauge manganin wires, whose low thermal conductance ensures a negligible heat leak to the ^{3}He liquid. The whole system was sealed in a vacuum can made of Stycast 1266 [16]. The vacuum can including the copper cold sink was immersed in the ^{3}He liquid.

The TEP S_{xx} is defined by $\nabla V_{xx} = S_{xx}\nabla T$, with the quantities measured in the following manner. A low frequency ($f_0 = 2.7$ Hz) ac voltage was applied to the heater and the T_1, T_2 (see inset, Fig. 1) were measured using the RuO_2 sensors and an ac bridge. The V_{xx} induced by the thermogradient was measured by a lock-in method at the frequency of $2f_0 = 5.4$ Hz. Both the T and voltage gradients were calculated using the dimensions given by the specimen.

In Fig. 1(a), we show the low-field magneto-TEP measured at 800 mK, for three electron densities. These traces reveal strong new oscillations appearing in $B < 3$ kG, where SdH oscillations are relatively weak. For example, up to four maxima can be clearly seen in the trace of $n_e = 2.03$. The arrows close to the traces indicate the maxima (indexed as $l = 1, 2, 3, 4$). A second derivative with

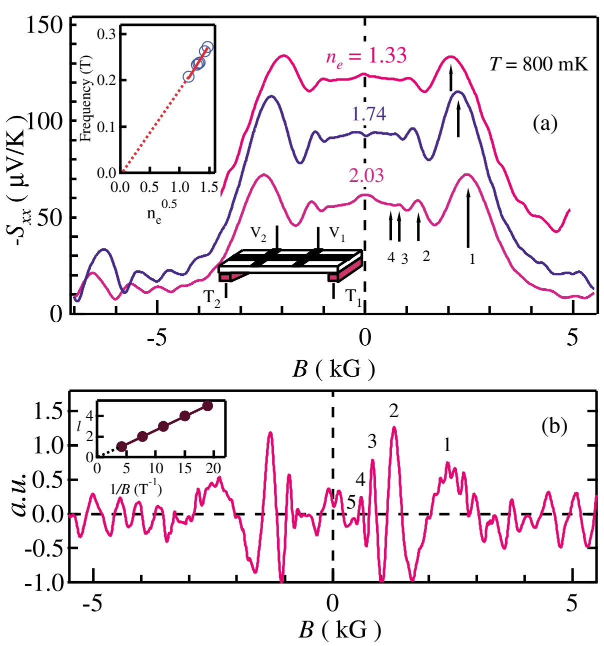

FIG. 1 (color online). (a) $-S_{xx}$ traces are shown for three densities n_e of 1.33, 1.74, and 2.03 in units of 10^{11} cm^{-2}, respectively; arrows indicate the maxima for $l = 1, 2, 3$, and 4 and the shift of the primary ($l = 1$) peak with increasing n_e. In (b), the second derivative against B for the high-density trace is shown; the numbers mark the oscillation peaks. The inset shows that oscillations are periodic in $1/B$.

respect to B, $-d^2S_{xx}/dB^2$, for the $n_e = 2.03$ is plotted in 1(b), together with a fan diagram (inset) showing the linear relation between the order l and the inverse B. We conclude from all three traces that the TEP oscillations are periodical in $1/B$. Moreover, the B position of the first peak scales with $\sqrt{n_e}$, as shown in the inset of Fig. 1(a). This behavior is distinct from that of SdH which scales with n_e, but consistent with the characteristics of a class [14,15] of weak B oscillations originating from cyclotron resonance with a $2k_F$ momentum transfer.

The TEP oscillations exhibit a remarkable T dependence, which, as will be shown later, can be attributed to inter-LL scattering by acoustic phonons. As an example, Fig. 2 shows the TEP data for $n_e = 2.03$ in a T range from 420 to 800 mK. Note that the TEP oscillation can be discerned at T as low as 300 mK for our experiment. With increasing T, the TEP signal increases dramatically, both at $B = 0$ and at $B < 5$ kG. At $B = 0$, the TEP S_0 is well understood as due to phonon drag. As will be shown, $S_0(T)$ is power law dependent in this experiment. As T rises, the TEP oscillation amplitude increases faster than S_0. This fact indicates that inter-LL scattering is dominant in the TEP oscillations observed here, since the inter-LL scattering mechanism leads to an exponential rather than a power-law T dependence. Indeed, at higher temperatures, $T \geq 800$ mK, the TEP oscillations become the dominating feature in TEP measurement at low field. As the cyclotron energy increases, the inter-LL scattering regime with an exponential-law T dependence disappears and the intra-LL-scattering SdH regime with a power-law T dependence prevails [12].

The nonoscillatory part of the TEP at $B < 3$ kG smoothly connects to TEP at $B = 0$. Our $B = 0$ TEP is consistent with previous results [8] with a similar phonon mean free path $\Lambda = 2.8$ mm, estimated from the thermal

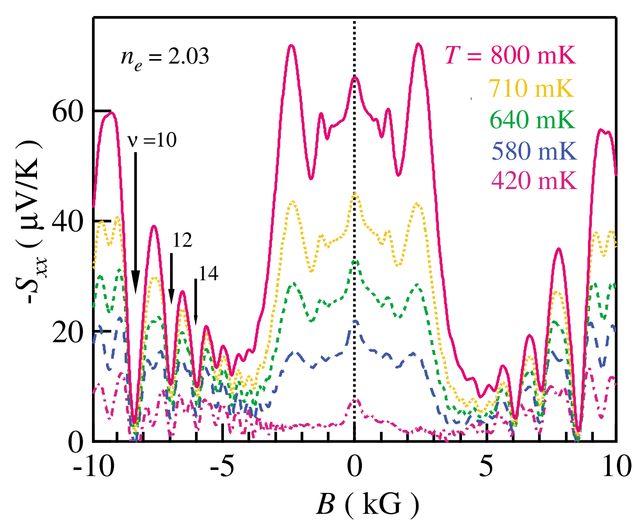

FIG. 2 (color online). $-S_{xx}$ (for $n_e = 2.03 \times 10^{11}$ cm^{-2}) traces at different temperatures show that the low field oscillations become stronger as T increases.

conductivity. This Λ is larger than the smallest dimension of the samples because of the multiple reflection at smooth surfaces. The precipitous drop of TEP beyond $\sim$3 kG is a generic feature for all of our samples, although quantitatively the ratio varies from sample to sample. The drop marks a transition from inter-LL to intra-LL scattering by phonons.

We now turn to a quantitative analysis for S_0 and the TEP oscillations. In Fig. 3(a), we plot S_0 vs T for three densities. All data show a power-law dependence on T, with an exponent between 3 and 4. This observation confirms quantitatively that the TEP at $B = 0$ is dominated by the phonon-drag mechanism [8]. The T-dependent amplitude of the first peak ($l = 1$) is presented in Fig. 3(b). The data strongly deviate from a power law, but can be fitted by a modified exponential relation $S_{xx} \sim \exp(-E/k_BT)/T^2$, predicted by Eq. (1). From the slope of the fit we arrive at the activation energies $E \sim 3.2$, 3.5, and 4.0 K, respectively, for the three densities, which are somewhat smaller than $\hbar\omega_c \sim$ 4.0, 4.4, and 4.8 K due to the LL broadening. These data strongly support the interpretation of a resonant phonon-drag mechanism being the origin of the TEP oscillations.

In principle, both 3D and 2D phonons can contribute to the resonant phonon-drag mechanism responsible for the oscillations. In the following, we start by considering the 3D case on more general ground. Qualitatively, the 2D case can be deduced from the 3D case.

We begin with the formula for the phonon-drag TEP in a magnetic field derived by Lyo [12]; a similar formulism was given by Kubakaddi *et al.* [11]. The TEP is obtained for a unit volume:

$$S_{xx} = \frac{-k_B h}{e\nu(k_BT)^2}\sum_{s\mathbf{q}}\sum_{n,n'} u_s\Lambda_{s\mathbf{q}}q_y^2 n_{s\mathbf{q}}|V_{s\mathbf{q}}|^2\Delta_z(q_z)\Delta_{n,\ell}(q_\parallel)\int d\varepsilon\rho_n(\varepsilon)\int d\varepsilon'\rho_{n'}(\varepsilon')f(\varepsilon)[1-f(\varepsilon')]\delta(\varepsilon+\hbar\omega_{s\mathbf{q}}-\varepsilon'), \quad (1)$$

where $\nu = \pi n_e l_B^2$ is the filling factor for spin-degenerate LLs, $n_{s\mathbf{q}}$ the Boson function for the phonon energy $\hbar\omega_{s\mathbf{q}}$, $f(\varepsilon)$ the Fermi function, u_s the sound velocity for the mode s, ρ_n the spectral function for the LLs, and $\Delta_z(q_z)$ is the conservation factor for q_z [12]. The square of the absolute value of the electron-phonon matrix element $|V_{s\mathbf{q}}|^2$ [12] is accompanied by the in-plane-momentum conservation factor:

$$\Delta_{n,\ell}(q_\parallel) = \frac{n!}{(n+\ell)!}\chi^\ell e^{-\chi}[L_n^\ell(\chi)]^2, \quad \chi = \frac{(q_\parallel l_B)^2}{2}, \quad (2)$$

which has a sharp principal maximum near $\chi = 4n$, namely, near the in-plane momentum transfer $q_\parallel \simeq 2k_F$ for $n \gg 1$, ℓ in view of $\varepsilon_F \simeq n\hbar\omega_c$ [14,17]. In Eq. (2), $L_n^\ell(\chi)$ is the associated Laguerre polynomial and $n' = n + \ell$ is the larger of n and n'. The phonon occupation factor restricts q_z to small values $q_z \ll 2k_F$ for this resonance at low temperatures. There are other secondary peaks below the main peak for $\Delta_{n,\ell}(q_\parallel)$. To reduce the computing time, we approximate the spectral density function $\rho_n(\varepsilon)$ by a rectangular distribution with a full width 2Γ centered at the LL energy $\varepsilon_n = \hbar\omega_c(n + 1/2)$ and take $\Gamma = 0.2$ meV for numerical evaluation.

The calculated TEP is plotted as a function of B in Fig. 4 for $\Lambda_{s\mathbf{q}} = 2$ mm for the three densities 2.03, 1.74, and 1.33×10^{11} cm^{-2} employing field-free screening for the electron-phonon interaction. Other parameters are

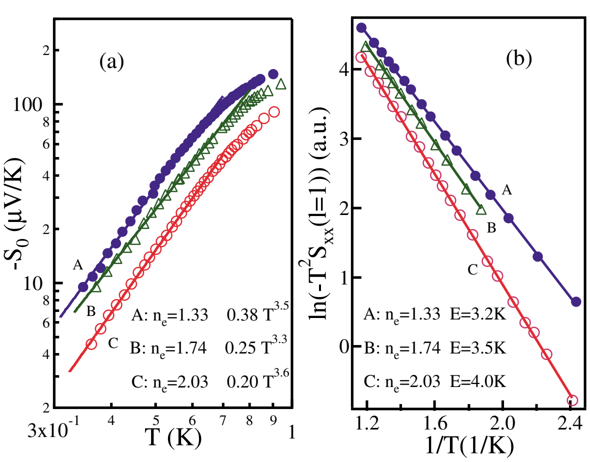

FIG. 3 (color online). (a) T dependence of the TEP at zero field (S_0) for three densities. The solid lines are fits to different power-law T dependences. (b) T dependence of the TEP at primary oscillation maxima ($l = 1$) for the three corresponding densities. Here, E is the fitted value for the phonon energy.

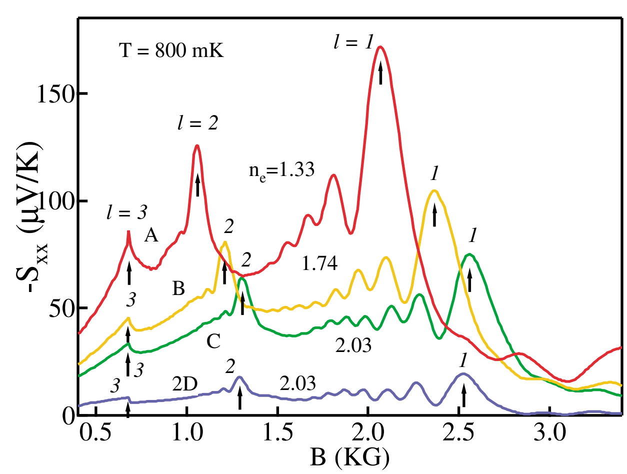

FIG. 4 (color online). Calculation of the TEP with a 3D-phonon model (upper three curves) and an estimate from a 2D-phonon model (bottom curve). The $l = 1$ maxima obey closely the T dependence in Fig. 3(b). With increasing n_e, the peak positions shift to higher B while the peak heights decrease.

well known and are given in Ref. [12]. The TEP is proportional to the $\Lambda_{s\mathbf{q}}$, which is basically an adjustable parameter. It is seen that the $\ell = 1, 2, 3, \ldots$ peaks are accompanied by secondary oscillations which are from the phonons with $q_{\parallel} < 2k_F$. Larger Γ makes the peaks broader and the valleys shallower compared to the sharper structures obtained for $\rho_n(\varepsilon) = \delta(\varepsilon - \varepsilon_n)$. *A*, *B*, and *C* traces display the TEP as a function of B for three densities $n_e = 1.33, 1.74$, and 2.03×10^{11} cm^{-2} at 0.8 K. It is seen that there is an approximate scaling relationship between the peak positions of B, satisfying $B \propto \sqrt{n_e}/\ell$. This relationship is consistent with the inter-LL resonance phonon picture $\hbar\omega_{2k_F} \simeq \ell\hbar eB/m^*$ in view of $\omega_{2k_F} \propto k_{\mathrm{F}} \propto \sqrt{n_e}$, yielding reasonable agreement with experimental data. We also find $S_{xx} \propto \exp(-E/k_BT)/T^2$ in agreement with Fig. 3(b) with $E \propto \sqrt{n_e}$ close to the transverse $2k_F$ phonon energy. Transverse phonons yield a dominant contribution (~70%) through strong piezoelectric scattering at low temperatures. Comparing with the data in Fig. 2, the calculated background TEP is much lower than the peaks, probably due to the simplistic non-self-consistent density of states employed in the present low-B situation where the LLs are closely separated. Also, the magnitude of the calculated TEP keeps decreasing as B approaches $B = 0$ in contrast to the data. At very low B, the number of the LLs become very large (e.g., $n > 100$ for $B < 0.4$ kG), requiring a zero-B formalism for a more accurate result.

The 2D phonon modes relevant in the GaAs/$\mathrm{Al}_x\mathrm{Ga}_{1-x}\mathrm{As}$ heterostructures are the leaky phonons [14,18] with phonon wave vector components $q_z = 0$ and $q = q_{\parallel}$. The TEP is again given by Eq. (1) with $\Delta(q_z) \equiv 1$ with the summation on q_z replaced by the summation over the leaky modes. For a rough estimate, we take same $|V_{s\mathbf{q}}|$ with the effective sample volume given by $\Omega = S\ell_p$, where S is the cross section of the well and ℓ_p is roughly the penetration depth of the mode. The result for the 2D phonons is compared with that of the 3D phonons in Fig. 4 for $\ell_p = 200$ Å using a pair of longitudinal and transverse modes. The B dependence of S_{xx} from the 2D phonons in Fig. 4 is very similar to that from the 3D phonons except that it is slightly shifted to lower B. The similarity between the result from the 2D and 3D phonons is numerically confirmed to hold at higher temperatures (e.g., 4 K). Since the magnetoresistance oscillations studied by Zudov *et al.* earlier originate from the same mechanism of phonon resonance [14], we believe that the result observed there at 4 K can also be caused by 3D phonons.

In Fig. 3(a), the linear contribution $S_0 \propto T$ from the electron-diffusion TEP is negligible at $B = 0$. This situation is similar to the data of Tieke *et al.* [8], but different from the results of Ruf *et al.* [6], in which the diffusive TEP is visible at 0.6 K and the data deviate from the T^3 slope at 0.6 K.

A high-mobility sample is necessary to observe this low-field acoustic-phonon resonance. In high-mobility samples, high LLs are not easily smeared out by impurity scattering. We expect stronger and sharper $2k_F$ oscillations for higher-mobility samples.

In conclusion, we have reported for the first time an oscillatory TEP in a weak magnetic field, where inter-LL scattering is accessible by acoustic phonons. The observation of such oscillations confirms a generic $2k_F$ momentum selection rule in electronic transport in a weak magnetic field where many LLs are occupied. Finally, it is found that while both 3D and 2D phonons cause qualitatively similar TEP oscillations in this regime, 3D phonons yield a substantial contribution to the oscillations to explain the data.

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