Experimental studies of the electron–phonon interaction in InGaAs quantum wires

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Electron-heating measurements are used to compare the form of the electron–phonon interaction in two-dimensional, and quasi-one-dimensional, InGaAs quantum wires. Evidence for a strongly *enhanced* interaction is found in the quasi-one-dimensional wire, and is suggested to result from the presence of the singularities in its electronic density of states. The Bloch–Gruneisen criterion is easily violated in this wire, and its energy-loss function is found to show a weak temperature dependence, which is argued to result from a saturation of scattering processes in the uppermost one-dimensional subband. [DOI: 10.1063/1.1495089]

The electron–phonon interaction (EPI) provides the dominant means for energy exchange in electron transport, and is crucially important to the operation of hot-electron devices.[1] In polar semiconductors, the interaction at low temperatures may involve the deformation potential and the piezoelectric interaction, both of which are influenced by electron screening. In nanostructures realized from these materials, these interactions may be further modified by the resulting confinement of electron (or phonon) motion. In this letter, we use electron-heating measurements to compare the form of the EPI in two-dimensional and quasi-one-dimensional InGaAs quantum wires. Evidence for a strongly *enhanced* interaction is found in the quasi-one-dimensional wire, and is suggested to result from the presence of the singularities in its electronic density of states.

A wire of width 600 nm and length 4 μm was formed[2] by photolithography and wet etching in an $In_{0.53}Ga_{0.47}As/In_{0.52}Al_{0.48}As$ delta-doped quantum well, grown on semi-insulating InP. This *quasi-two-dimensional* wire (see the discussion below) was employed as a reference structure, to compare with the heating characteristics of a much narrower (25 nm), trench-type, wire, whose fabrication and equilibrium-transport properties have been described previously.[3] Both wires were mounted in contact with the mixing chamber of a dilution refrigerator and measurements of their magnetoresistance were made using standard lockin (11.5 Hz) techniques. The measurements here were performed at lattice temperatures (T_0) in the range of 0.01 to 10 K, using drive currents that varied from 0.1 to 100 nA.

Fig. 1 shows the magnetoresistance of the 600 nm wire at a set of different lattice temperatures. Shubnikov–de Haas

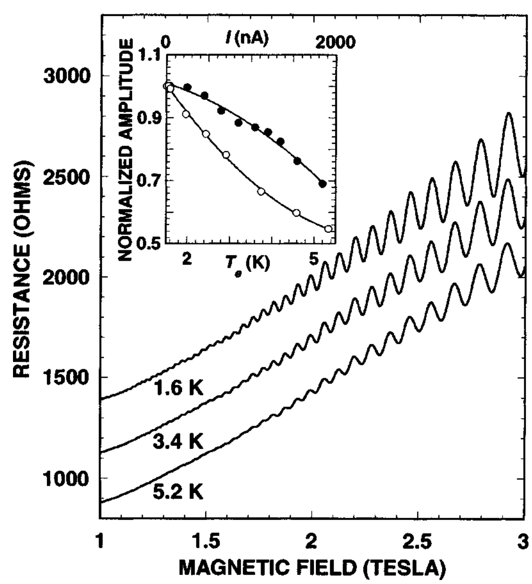

FIG. 1. The magneto-resistance of the 600 nm wire at three different lattice temperatures (T_0). Traces at 3.4 and 5.2 K are offset by −250 and −500 Ω, respectively. A constant current of 5 nA was used in all three measurements. Inset: dependence of the Shubnikov–de Haas oscillation amplitude on lattice temperature (filled circles) and current (open circles). The amplitude is normalized to the value obtained at the lowest temperature, or current, respectively, and the data points are an average over the range from 2.5 to 4.0 T. Solid lines are guides to the eye.

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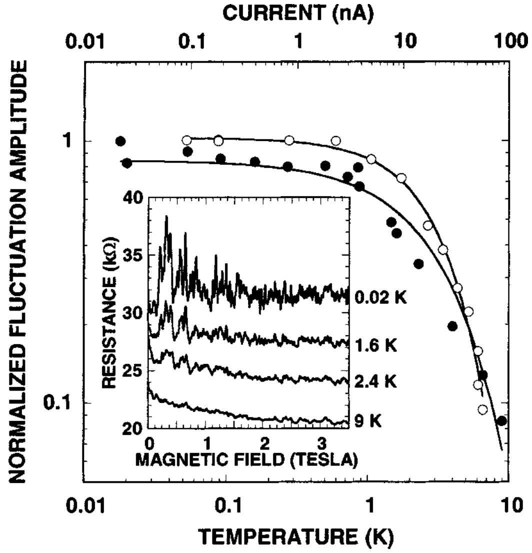

FIG. 2. Main panel: dependence of the conductance fluctuation amplitude on lattice temperature (filled circles) and current (open circles) in the 25 nm wire. Inset: temperature dependence of the magnetoresistance of the trench-type wire. The measurement current was 0.1 nA and the curves at 2.4, 1.6, and 9 K are offset by +3, +6, and +9 kΩ, respectively.

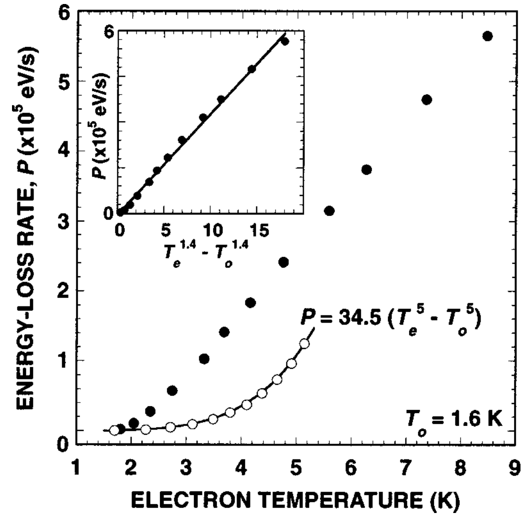

FIG. 3. Variation of the energy-loss rate with temperature in the 600 nm (open circles) and 25 nm (filled circles) wires. In the inset we replot the data for the 25 nm wire to show the $T^{1.4}$ variation of $F(T)$.

oscillations are clearly seen at higher magnetic fields and are used to infer a carrier density $n_s = 2.6\times 10^{12}\ \mathrm{cm}^{-2}$. The oscillations show *no* deviation from inverse-field periodicity at low magnetic fields,[4] *indicating that the quasi-one-dimensional levels of this wire are energetically obscured*. At temperatures below ~3 K, the Shubnikov–de Haas oscillations onset at a temperature-independent magnetic field of ~1 T. Equating this with the point at which the cyclotron orbit begins to fit inside the wire, we estimate an effective width of ~500 nm, suggesting that the wet etching of this wire gives rise to a total sidewall depletion of 100 nm. From the zero-field resistance of the wire (~1260 Ω), we estimate a wire mobility of 15300 cm^2/Vs.

To study electron heating in the 600 nm wire, we use the Shubnikov–de Haas effect as an electron thermometer. An increase of the drive current at fixed lattice temperature produces a decrease in the oscillation amplitude, similar to the effect of increasing the lattice temperature. By comparing the current-induced change in the oscillation amplitude to that which arises when the temperature is varied with only a small drive current (Fig. 1, inset), we can infer the current dependence of the electron temperature (T_e) for this wire. We return to a discussion of this data shortly.

In contrast to the above behavior, the magnetoresistance of the 25 nm wire is dominated by reproducible fluctuations[3] (Fig. 2, inset), which result from a quantum interference effect.[1] Due to its extremely narrow width, no evidence for the Shubnikov–de Haas effect is found in this wire.[3] However, it is known that we may use the fluctuation *amplitude* as an electron thermometer.[5] In the main panel of Fig. 2, we plot the variation of this amplitude with temperature and drive current, and this data allows us to infer the current dependence of the electron temperature in this wire. (The current variation was determined at a cryostat temperature of 1.6 K. As can be seen from the main panel of Fig. 2, the fluctuation amplitude does not increase significantly below ~1 K.)

In Fig. 3, we summarize the results of our heating studies of the two wires by plotting the *energy-loss rate per electron in the wire* as a function of the electron temperature. The energy-loss rate is simply I^2R/n_sWL, where I is the drive current, n_s is the carrier density, and R, W, and L are the average resistance, the width, and the length, of the wire, respectively. The open symbols in Fig. 3 are for the 600 nm wire, while the solid symbols are for the 25 nm wire. Since the values of n_s and W cannot be determined directly for this wire, we instead *assume* $n_s = 3\times 10^{12}\ \mathrm{cm}^{-2}$ and $W = 25$ nm, which our recent studies of InGaAs *ridge-type* quantum wires[4] suggest should be reasonable estimates. For the purposes of the discussion below, however, we note that these choices for n_s and W will *not* affect the *functional* dependence of the energy-loss rate on T_e.

The data of Fig. 3 suggest a *modified* EPI in the 25 nm wire, and our discussion is based on an analysis of the *energy-loss function*. This is expected to exhibit the polynomial form $F(T) = AT^n$, where the constant A and index n reflect the details of the relevant EPI mechanism.[6–8] The electron *energy-loss rate* is then given by

$$P(T) = F(T_e) - F(T_0). \quad (1)$$

We begin by considering the behavior exhibited by the 600 nm, two-dimensional, wire. The line through the open symbols in Fig. 3 is a fit to the form of AT^n, with $A = 34.5\ \mathrm{eV\,s^{-1}K^{-5}}$ and $n = 5$. For the temperature range studied here, we expect that the Bloch–Gruneisen criterion should be satisfied for the 600 nm wire. The condition for this limit may be expressed as $q_T \ll 2k_F$, where $q_T = 2.8k_BT/\hbar u$ is the wave vector of a thermal phonon and u is the sound velocity, indicating that the EPI predominantly involves scattering events with *small* momentum transfer. (Due to a lack of experimental data for InGaAs, we assume that u is in the range of 3 to $5\times 10^3\ \mathrm{m\,s^{-1}}$.) The investigated temperature range also corresponds to the limit of *weak* screening, the crossover to which is defined by the condition $q_T\lambda_s \approx 1$. Assuming a Bohr radius (λ_s) for InGaAs of 9.5 nm, weak screening is expected for temperatures above ~1 K. We therefore suggest that the observed T^5 variation in the

600 nm wire results from electron scattering from the *unscreened deformation potential*. The energy-loss function due to this mechanism is predicted to take the form:[6,7]

$$F_{u-d}=A_{u-d}T^5=\frac{6\zeta(5)k_B^5}{\sqrt{2\pi^5}\hbar^7}\frac{m^{*2}U^2}{n_3^{3/2}\rho u^4}T^5. \tag{2}$$

Here, ζ is the zeta function, ρ is the mass density $(=5.4\ \mathrm{g\,cm^{-3}})$,[9] and U is the deformation potential. For InGaAs, we take $U=16\ \mathrm{eV}$,[10] and assume a sound velocity $u=3.3\times10^3\ \mathrm{m\,s^{-1}}$ (this latter value is that for InGaP, whose acoustic parameters are similar to those of InGaAs[10]). In this way, we find a prefactor $A_{u-d}=30\ \mathrm{eV\,s^{-1}K^{-5}}$ from Eq. (2), which is in agreement with our findings for the 600 nm wire (see Fig. 3). For completeness, however, we should point out that the influence of nonparabolic energy bands is not included in Eq. (2), and may be important for the energy-loss rate in InGaAs.

A comparison of the data in Fig. 3 reveals that, over the range of temperature studied, the energy-loss rate is larger in the narrow wire, consistent with an enhanced EPI in this structure. Such an enhancement has actually been predicted for narrow wires, due to the singularities in the quasi-one-dimensional density of state.[11] Since the occupation of electron states in the lower subbands is expected to be strongly degenerate, the details of electron–phonon scattering should be dominated by the contribution of electrons in the highest subband.[11] This point is important for explaining the weak power-law variation of $F(T)$, observed in the narrow wire. Also important is the fact that the Bloch–Gruneisen limit for this wire should be restricted to *much lower* temperatures than those studied experiment. A constraint on small-angle scattering does arise at temperatures greater than $hu/2.8k_BW$, where the one-dimensional subband structure of the wire limits the momentum transfer to of order $1/W$.[11,12] Also, electron–phonon scattering in the highest subband may involve *backscattering* processes, *even at low temperatures*, and this mechanism for energy transfer to the phonon subsystem is known to be more effective than processes involving only small-angle scattering.[13] Once this occurs, the electron–phonon matrix element and the phonon phase-space volume available for electron–phonon scattering become independent of temperature and the electron–phonon relaxation time ($\tau_{\mathrm{e-ph}}$) should saturate. Since the energy-loss function may be expressed as $F(T)=C_eT/n_sn\tau_{\mathrm{e-ph}}$[14] (where C_e is the electron heat capacity and should depend only weakly on temperature for nondegenerate filling of the uppermost subband), a weak power-law variation of $F(T)$ is expected in this regime. Indeed, in the inset to Fig. 3, we show that $F(T)$ in the 25 nm wire is well described by a $T^{1.4}$ dependence, implying that $C_e\propto T^{0.4}$, consistent with the notion of an EPI interaction that is dominated by electron scattering in the weakly degenerate uppermost subband.

In conclusion, we have used electron-heating measurements to study the form of the EPI in two-dimensional, and quasi-one-dimensional, InGaAs quantum wires. Evidence for a strongly *enhanced* EPI was found in the narrower wire and has been suggested to result from the strong backscattering of electrons due to the singularities in the one-dimensional electronic density of states. An energy loss function with a weak power-law index ($n\sim1.4$) was also found for this wire and has been argued to be consistent with a saturation of electron backscattering processes in the highest occupied one-dimensional subband.

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