Radiation of acoustic phonons from quantum wires

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We have investigated by the Monte Carlo technique the radiation of ballistic acoustic phonons from quasi-one-dimensional electron gas in quantum wires. At low temperatures and over a wide range of electric fields, all excess heat in quantum wires is dissipated by means of acoustic phonons. Due to the uncertainty of momentum conservation during electron–acoustic-phonon scattering, electrons emit acoustic phonons with large transverse momentum components. Consequently, in this transport regime quantum wires radiate fluxes of nonequilibrium acoustic phonons into surrounding material. Nonequilibrium acoustic phonons can propagate ballistically over macroscopic distances. Ballistic fluxes of nonequilibrium acoustic phonons have been previously detected experimentally in quantum well structures. We have calculated the angular and energy spectrum of nonequilibrium acoustic phonons radiated from quantum wires. © *1995 American Institute of Physics.*

INTRODUCTION

Dynamics of nonequilibrium acoustic phonons in nanostructures has recently attracted significant attention of the scientific community.[1–15] If the lattice temperature is low so that the phonon–phonon interactions are negligible, the acoustic phonons with frequencies up to 0.8 THz (acoustic-phonon energy up to 3 meV) propagate ballistically over macroscopic distances. Radiation of ballistic acoustic phonons from the slightly heated quasi-two-dimensional (2D) electron gas has been experimentally detected[1–3] using sensitive thin film bolometers (particularly, optically activated bolometers) or superconducting tunnel junctions. The ballistic flux of nonequilibrium acoustic phonons not only carries information about the electron system which generated this flux but also provides a very efficient means of excess heat removal from nanoscale devices.

The previously referenced articles have dealt primarily with the kinetics of acoustic phonons, their scattering, propagation, and decay. It must be noted, however, that the nature of acoustic-phonon emission itself by a low dimensional electron gas is substantially different from that in bulk materials. The peculiarities of electron interactions with acoustic phonons in nanostructures are reflected in the character of nonequilibrium acoustic-phonon propagation in real space. In this article, we want to address the problem of acoustic-phonon emission by heated electrons, i.e., we will consider generation of nonequilibrium acoustic phonons and will determine how this generation defines acoustic-phonon transport. This initial stage of acoustic phonon kinetics, namely, emission of nonequilibrium acoustic phonons by heated low dimensional electron gases, is of crucial importance since it defines the efficiency of excess heat dissipation in nanostructures and nanodevices. The radiation of acoustic phonons should be even more pronounced from quasi-one-dimensional (1D) quantum wires (QWIs), where at low temperatures in a wide range of electric fields acoustic-phonon emission is the sole electron scattering mechanism.[16] However, no experimental or theoretical studies of nonequilibrium acoustic-phonon generation and radiation from QWIs have yet been reported. We report here the first results on the Monte Carlo simulation of nonequilibrium acoustic-phonon emission by a heated electron gas in GaAs/AlGaAs QWIs. We have calculated angular and energy distributions of acoustic phonons and determined the preferable directions and nature of their propagation from a heated 1D electron gas.

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RESULTS AND DISCUSSION

We will be interested in the electron transport regime controlled by acoustic-phonon emission,[16] because in this regime we may expect significant generation of nonequilibrium acoustic-phonon populations. This transport regime is possible if $\hbar u \Delta q \gg k_B T$, where $\hbar u \Delta q$ is the acoustic phonon energy defined by the uncertainty $\Delta q = 2\pi/L$ of conservation of the transverse component of phonon wave vector, u is sound velocity, L is the effective thickness of the QWI $L^{-2} = L_y^{-2} + L_z^{-2}$.[17] The above condition is fulfilled for $L \leqslant 100$ Å and the lattice temperature $T = 4$ K or less. We will consider a QWI with a cross section of 80×80 Å^2 at an equilibrium lattice temperature of $T = 4$ K. We consider a nondegenerate electron gas. Our model is based on an infinitely deep potential well for electrons. We have taken into account multisubband energy structure and electron scattering by confined longitudinal optical (LO), localized interface (surface) optical (SO) phonons,[18] as well as by bulk-like acoustic phonons.[17] Although in general, acoustic phonons in nanostructures may have several modes which differ from bulk modes,[19–21] the model of bulk-like acoustic phonons is justified, if materials constituting a nanostructure have similar elastic properties as for the case of GaAs and AlGaAs.

Let us first choose the electric field range, where acoustic-phonon emission is most effective. Figure 1 demonstrates the relative efficiency of various scattering mechanisms versus electric field. One can see that in the range of

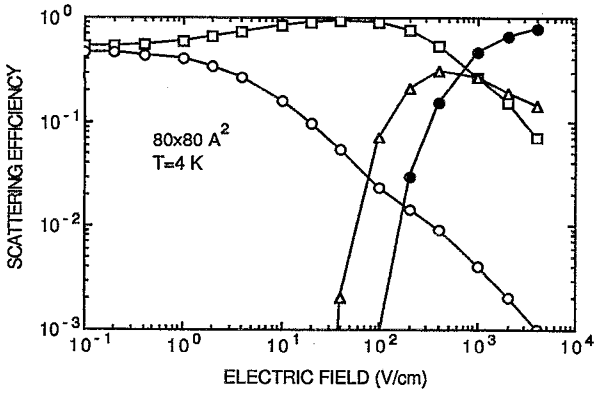

FIG. 1. Relative efficiency of different scattering mechanisms vs electric field. Open squares represent efficiency of acoustic-phonon emission, open circles—acoustic-phonon absorption, open triangles represent efficiency of surface optical-phonon emission, and full circles—confined longitudinal optical-phonon emission. Quantum wire cross section is 80×80 Å^2, lattice temperature T=4 K.

electric fields of 2–200 V/cm the emission of acoustic phonons is the sole important mechanism, exceeding all other scattering rates by several orders of magnitude. At higher fields optical phonon emission comes into play. Thus, the total excess energy pumped into the electron system of the QWI in the field range of 2–200 V/cm is dissipated through the emission of acoustic phonons. Figure 2 shows the total Joule power, $P_{\text{Joule}}=ev_dE$, dissipated by a single average electron, where v_d is the electron drift velocity and the power dissipated via emission of acoustic phonons as a function of electric field. One can see from Fig. 2, that within the numerical error the total Joule power dissipated is equal to the power dissipated through acoustic phonons in the above mentioned field range. It is important to stress here that the power increases by five orders of magnitude when the electric field is increased from 0.1 to 100 V/cm. Consequently, the power carried out of the QWI by the acoustic-phonon flux exceeds the background radiation from the QWI

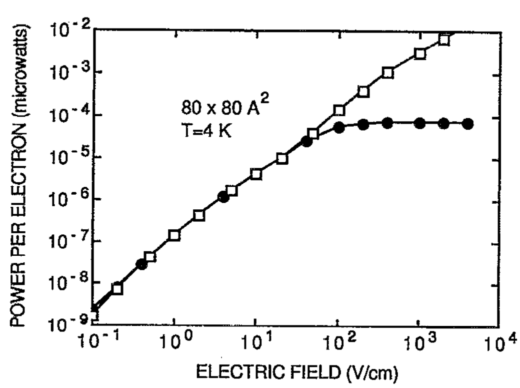

FIG. 2. Power emitted by a single electron via acoustic phonons (open squares) and the total Joule power dissipated per single electron, ev_dE, (full dots), vs applied electric field in 80×80 Å^2 quantum wire at T=4 K lattice temperature. Here, v_d stands for electron drift velocity and E for electric field.

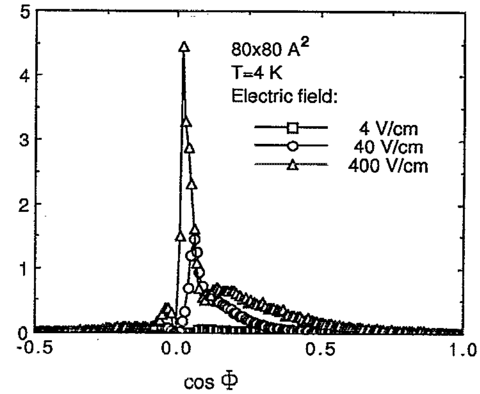

FIG. 3. Angular distribution of the power flux carried by nonequilibrium acoustic–phonons for three different electric fields. cos $\Phi=q_x/q$, where q_x is the along-wire component of acoustic-phonon wave vector, and q is the total acoustic-phonon wave number. Quantum wire cross section is 80×80 Å^2, lattice temperature T=4 K.

by five orders of magnitude. Therefore, acoustic-phonon fluxes from QWIs should be experimentally detectable.

The angular distribution of acoustic phonons emitted by electrons carries much information about the electron system. The angular distribution of acoustic phonons emitted by hot 2D electron gases has been studied experimentally.[2,3] It is obvious from the conservation laws that the perpendicular component of phonon wave vector emitted by a 1D electron gas is almost always greater than the longitudinal component. Indeed, from energy conservation it follows that

$$q_x^2\left[\left(\frac{\hbar}{2m^*u}\right)^2(q_x\mp2k_x)^2-1\right]=q_T^2,$$

where q_x and q_T stand for longitudinal and transverse components of phonon wave vector $\mathbf{q}$, and k_x denotes the electron wave number. The factor in the square brackets on the left-hand side of the above equation is much larger than 1 if

$$\frac{\hbar^2(q_x\mp2k_x)^2}{2m^*}\gg m^*u^2.$$

By substituting numerical values for GaAs one can show that the above inequality holds for forward scattering (i.e., scattering during which the sign of the electron momentum does not change) of electrons with energies higher than 0.01 meV. Due to the uncertainty of transverse momentum conservation during 1D electron interaction with bulk-like acoustic phonons the typical transverse component of wave vector of acoustic phonons emitted by a heated 1D electron gas is of the order of $\Delta q=2\pi/L$. Simple estimates show that the inequality $q_x\ll q_T\approx\Delta q$ also holds for backward scattering (i.e., scattering during which the sign of the electron momentum changes) of electrons with energies higher than 0.01 meV and less than about 100 meV for an 80×80 Å^2 QWI. As a result, electrons in a wide range of energies emit acoustic phonons which propagate almost perpendicularly to the QWI.

Figure 3 shows the angular dependence of the power flux associated with ballistic acoustic phonons from QWIs calculated by the Monte Carlo technique. The sharp peak

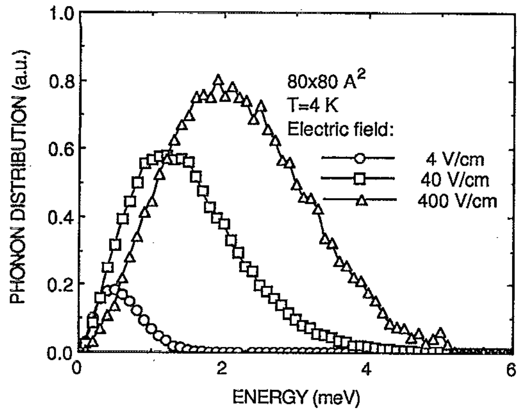

FIG. 4. Energy distribution of acoustic phonons emitted by electrons for three different electric fields. Quantum wire cross section is 80×80 Å^2, lattice temperature T=4 K.

close to cos $\Phi = q_x/q = 0$ proves that acoustic phonons are radiated predominantly in a direction perpendicular to the QWI. At higher electric fields a second lower and broader maximum appears at higher positive values of cos Φ. This maximum is caused by backward electron scattering with phonon emission. Backward scattering always involves greater values of the longitudinal component of phonon wave vector q_x. Consequently, these phonons propagate at smaller angles to the QWI. When increasing the electric field the mean electron energy increases. So does the q_x component of acoustic phonons emitted during backward scattering. Therefore, the maximum of the angular distribution related to backward scattering moves toward a higher magnitude of cos Φ and separates from the sharp peak related to forward scattering. Hence, the angular distribution of the acoustic phonon flux contains valuable information about the electron distribution function.

Figure 4 depicts the energy distribution of emitted acoustic phonons. The energy spectrum of emitted acoustic phonons determines the character of acoustic-phonon propagation. The phonon–phonon scattering rate is proportional to ω^4 and phonon-defect scattering to ω^5. Simple estimates show that the mean free paths of acoustic phonons with energies up to 5 meV are of the order of a millimeter or larger.[13] Therefore, acoustic phonons with energies less than about 5 meV propagate ballistically over macroscopic distances. As one can see from Fig. 4, all emitted acoustic phonons have energies in the range of 0–5 meV; i.e., all such phonons can propagate ballistically.

SUMMARY

Summarizing, we have calculated by the Monte Carlo technique the emission of acoustic phonons from the electron gas of a 1D quantum wire. Our results show that in a wide range of electric fields acoustic-phonon emission is so effective that it is virtually the sole scattering mechanism. The total power pumped into the electron system by the electric field in steady-state is dissipated through acoustic phonons. Our estimates show that electron-generated nonequilibrium acoustic phonons propagate ballistically and almost perpendicularly to a quantum wire. Fluxes of nonequilibrium acoustic phonons should be easily detectable experimentally since their radiation density exceeds the thermal equilibrium radiation by five orders of magnitude.

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