Electrophonon Resonances in Mesoscopic Structures

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A new type of resonant behavior is predicted for the hot electron 1D ballistic transport regime in nanostructures. The resonances are manifested by the sharp drops in current which occur when (i) the separation between a pair of levels of transverse quantization equals the energy of an optical phonon $\hbar\omega_0$, and (ii) a condition for generation of optical phonons $eV > \hbar\omega_0$ is met, where V is the voltage bias across the conductor.

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We investigate a new type of resonant behavior that can occur in quasi-1D (Q1D) nanostructures in the hot electron regime. The resonant condition is of the form

$$\hbar\omega_0 = \epsilon_m(0) - \epsilon_n(0). \tag{1}$$

Here ω_0 is the frequency of the long wavelength optical phonons, whereas $\epsilon_m(0)$ is the threshold of propagation in the mth Q1D subband (channel). In addition, the overall condition for the generation of the optical phonons $eV > \hbar\omega_0$ has to be met along with Eq. (1) in order for the resonances to take place [1]. Sharp resonances may be achieved only for sufficiently low temperatures.

The physics of this phenomenon is associated with singularities peculiar to the 1D electron density of states. The same kind of singularity is also typical for a 3D electron spectrum in a magnetic field. Hence there is a close resemblance between this phenomenon and the magnetophonon resonance [2]. The term *electrophonon resonance* (EPR) was first introduced by Bryskin and Firsov [3] who have predicted EPR for nondegenerate semiconductors in a very strong electric field (cf. also with [4,5]). Although in our case the resonances are achieved by the variation of the electric voltage as well, the underlying physics is quite different. Such voltage can be either the bias voltage across the nanostructure or the gate voltage whose variation can result in a variation of the electron spectrum.

We begin with a description of the main features of the background on which the resonant structure is superimposed. In the zeroth approximation, the collisions can be neglected and the electron transport may be considered as ballistic. In the linear response regime (see [6]), the conductance G is a steplike function of the Fermi level or the gate voltage. Each step corresponds to inclusion of a new mode of transverse quantization to the conduction process. The height of each step is equal to the quantum of conductance, $G_0 = 2e^2/h$, and according to the Landauer formula is multiplied by a prefactor whose physical interpretation is that of a transmission probability [7]. It is unity for a uniform 1D conductor which we consider here. In such a case, the electron spectrum consists of a set of 1D subbands. Each is characterized by the dispersion law $\epsilon_n(p) = \epsilon_n(0) + p^2/2m^*$, where n is the index for the Q1D subbands, p is the x component of the electron quasimomentum, while m^* is the electron effective mass.

With an increasing bias voltage V, one could expect a non-Ohmic behavior of the conductance. A substantial deviation of current from the Ohmic value may be expected if the ratio eV/μ is not small—cf. with Ref. [8]. The functional dependence of $J^{(0)}(V)$ may be rather peculiar, so that the differential conductance $g = dJ^{(0)}(V)/dV$ appears to be an oscillating function of the potential difference V [9] (see Fig. 1). Another important parameter is k_BT. For $eV \ll k_BT$ the dependence $J^{(0)}$ on V is purely Ohmic, with an exponential accuracy. However, $g(V)$ exhibits oscillations for $eV \gtrsim k_BT$.

The distribution function of electrons, $f^{(0)}(\epsilon)$, in the subband n is given by $f^{(0)}(\epsilon) = f^{(F)}(\epsilon - \mu^{(\pm)})$, where $f^{(F)}$ is the Fermi function while $\epsilon = \epsilon_n(p)$ is the electron energy. The upper (lower) sign is for $p > 0$ ($p < 0$) that corresponds to electrons coming from the left (right) reservoir (cf. with Ref. [6]), and $\mu^{(\pm)} = \mu \pm eV/2$ is the corresponding chemical potential.

The current $J^{(0)}$ is given by [10]

$$J^{(0)} = \frac{e}{\pi\hbar\beta}\sum_n \ln\left[\frac{1 + \exp(\beta\mu_n^{(+)})}{1 + \exp(\beta\mu_n^{(-)})}\right], \tag{2}$$

where $\mu_n^{(\pm)} = \mu^{(\pm)} - \epsilon_n(0)$ and $\beta = 1/k_BT$. Here we assume that the chemical potential μ is determined by the chemical potentials of the reservoirs with which the ballistic conductor is in contact. Half steps of the differential conductance in basic units of G_0 can be seen in

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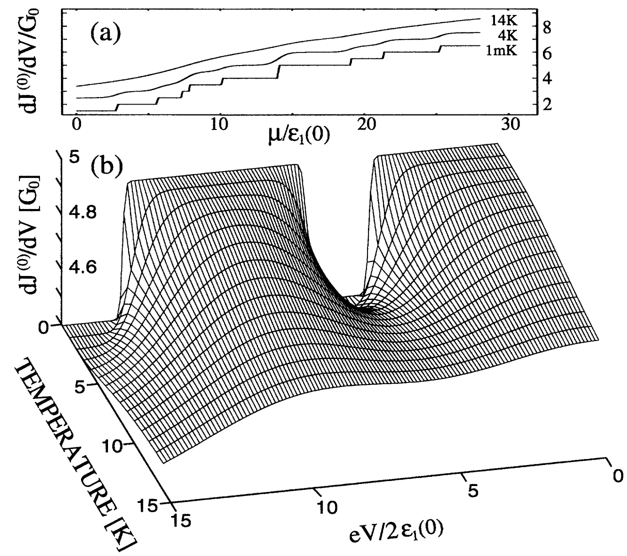

FIG. 1. $dJ^{(0)}/dV$ as a function of temperature and (a) chemical potential, μ [at $eV/\epsilon_1(0) = 10$], or (b) voltage bias, eV. Curves are vertically offset by one unit in (a). A constant chemical potential of 14 meV and $L_y = 100$ nm [$\epsilon_1(0) = 0.56$ meV] are assumed.

Fig. 1 and have been predicted by Glazman and Khaetskii [11]. They focus on the effects of nonadiabaticity of the constriction geometry as a cause of smearing of quantum steps. Here we wish to emphasize an important role of *temperature-sensitive features* of the nonlinear behavior given by Eq. (2). The first experimental evidence of these additional plateaus has been found by Kouwenhoven *et al.* [8]; further important experimental evidence was obtained by Patel *et al.* [12].

In order to calculate the variation of the total current $\Delta J = J - J^{(0)}$ due to the electron-phonon interaction, we assume that $|\Delta J|/J^{(0)}$ is small and make use of the perturbative methods. We are particularly interested in the regions where the total current J drops sharply. These features are especially pronounced at low temperatures and are due to the resonant nature of the intersubband electron-optical phonon scattering.

The total EPR contribution to the current is a sum of currents carried by individual channels [13]. In turn, a *channel current* ΔJ_n is made up of *partial currents* $\Delta J_{nn'}$ due to particular interchannel transitions:

$$\Delta J = \sum_n \Delta J_n = \sum_{nn'} \Delta J_{nn'} . \tag{3}$$

The partial currents are given by

$$\Delta J_{nn'} = -2eL_x(1 - e^{-\beta eV}) \int \frac{dp\, dp'\, d^d q_\perp}{(2\pi)^{d+2}\hbar^2}\, \theta(-pp') \times C_{nn'}(\mathbf{q})(\mathcal{A}_{nn'}^{(-)} + \mathcal{A}_{nn'}^{(+)})\delta(\epsilon' - \epsilon - \hbar\omega_{\mathbf{q}}) . \tag{4}$$

Here $\epsilon = \epsilon_n(p)$, $\epsilon' = \epsilon_{n'}(p')$,

$$\mathcal{A}^{(+)} = F_\mu(\epsilon', \epsilon, V)(1 + N_{\mathbf{q}}),$$
$$F_\mu(\epsilon', \epsilon, V) = f^{(F)}(\epsilon' - \mu^{(+)})[1 - f^{(F)}(\epsilon - \mu^{(-)})],$$
$$C_{nn'}(\mathbf{q}) = W_{\mathbf{q}} |\langle n'| \exp(i\mathbf{q}_\perp \cdot \mathbf{r}_\perp) |n\rangle|^2,$$

where $\mathcal{A}^{(-)}$ is obtained from $\mathcal{A}^{(+)}$ by replacing ϵ by ϵ' and $1 + N$ by N; $d + 1$ is the dimension of the system, L_x is the total length of the conductor, $N_{\mathbf{q}}$ is the average phonon occupation number (which we assume to be the equilibrium Bose function), and $\langle \mathbf{r}_\perp | n\rangle$ is the transverse part of the electron wave function [14]. The term proportional to $\mathcal{A}^+$ describes the phonon emission, whereas $\mathcal{A}^{(-)}$ corresponds to absorption processes. The summation over all the phonon branches is implicit. As expected, ΔJ diminishes the total current.

We consider the scattering of electrons by the three-dimensional bulk polar optical phonons. Then $W_{\mathbf{q}} = 2\pi he E_0/m^* q^2$, where $eE_0 = m^* e^2 \omega_0(\varepsilon_0 - \varepsilon_\infty)/\hbar\varepsilon_\infty\varepsilon_0$; ε_∞ and ε_0 are the dielectric constants for infinite and zero frequencies, respectively [15]. We ignore the dispersion of optical phonons, $\omega_{\mathbf{q}} = \omega_0$.

In the low temperature limit [$N(\omega_0) \ll 1$], the nth channel contribution to the ΔJ is then

$$\Delta J_n = -\frac{em^* L_x}{4\pi^2\hbar^2}(1 - e^{-\beta eV}) \sum_m \int \frac{d^2 q_\perp}{(2\pi)^2} W_{\mathbf{q}} C_{nn'}(\mathbf{q}_\perp) \times \int_{R_{nm}}^{\infty} d\epsilon\, \frac{F_\mu(\epsilon + \omega_0, \epsilon, V)}{[\epsilon(\epsilon - R_{nm})]^{1/2}}\, \theta(\epsilon)\theta(\epsilon V - \hbar\omega_0), \tag{5}$$

where $R_{nm} = \epsilon_m(0) - \epsilon_n(0) - \hbar\omega_0$.

The singularities in (5) are due to the initial and final densities of states. As long as $R_{nm} \neq 0$, they are integrable, and the integral in (5) is well behaved as a function of bias or gate voltage. However, when the resonant condition of Eq. (1) is met ($R_{nm} \to 0$), the denominator in Eq. (5) amounts to a $1/\epsilon$ singularity which renders (5) nonintegrable (cf. with [2]); this leads to a logarithmic divergence [16] (which is, of course, removed if one takes into consideration scattering events of high order) and, consequently, to a significant increase in the resistance. Since no such singularities are present in 2D and 3D cases for an otherwise equivalent physical system, no EPR in our sense of the word is to be expected there either.

The integral in (5) vanishes unless the functions $f^{(F)}(\epsilon + \hbar\omega_0 - \mu^{(+)})$ and $1 - f^{(F)}(\epsilon - \mu^{(-)})$ overlap. The presence of such an overlap in $\mathcal{A}^{(+)}$ of Eq. (4) at low temperatures leads to the condition $eV > \hbar\omega_0$ (to within k_BT) and is explicitly taken into account by the step function $\theta(eV - \hbar\omega_0)$.

Let us investigate a resonant behavior of a partial current, say, ΔJ_{12}. In Fig. 2 we show schematically the transitions that correspond to ΔJ_{12}. If the position of the Fermi level μ is exactly halfway between the bottoms of the subbands 1 and 2 and the applied voltage bias slightly exceeds the optical energy [see Fig. 2(a)], the transition

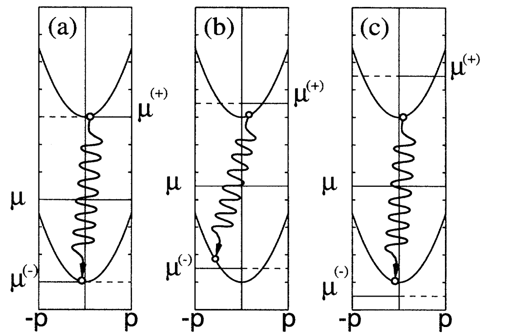

FIG. 2. While the conditions for EPR are favorable in (a) and (c), no EPR is possible for (b).

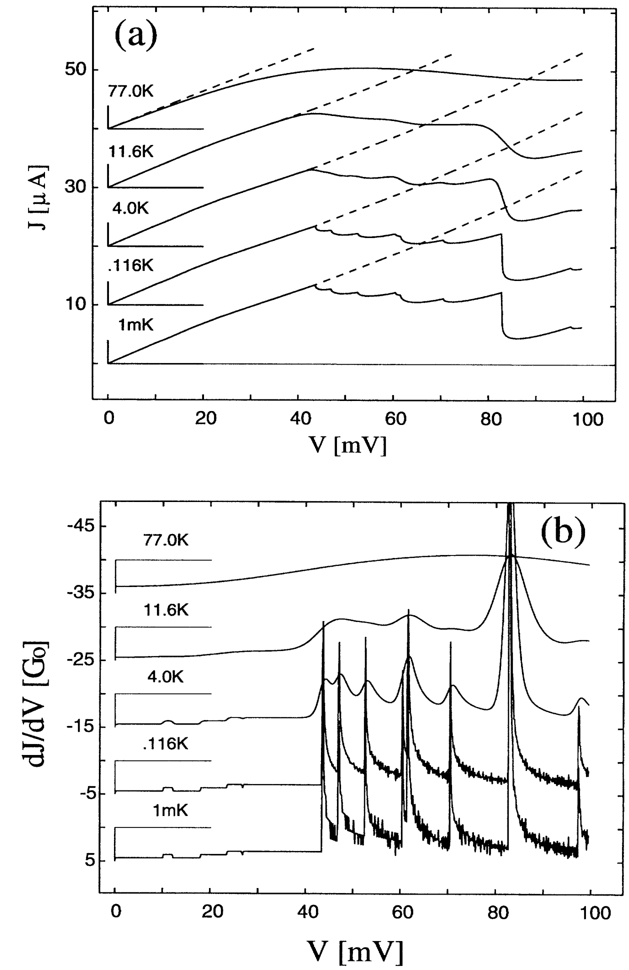

FIG. 3. (a) The current J and (b) differential conductance $g(V)$ as a function of V for various temperatures ($L_y = 100$ nm). Curves are vertically offset by 10 μA in (a) and by $10G_0$ in (b). A constant chemical potential of 14 meV is assumed.

indicated by the arrow is allowed. Such a transition would correspond to the electrophonon resonance.

Indeed, in this case, the right-hand part ($p > 0$) of band 2 near its bottom is occupied by electrons, whereas the left-hand part ($p < 0$) of band 1 is empty. Therefore, the transitions where an electron is backscattered by an optical phonon from the bottom of band 2 to the bottom of band 1 are allowed. If, on the other hand, the position of the Fermi level is shifted by, say, $\Delta\mu$ upward from the midpoint between the bottoms of subbands 1 and 2 [see Fig. 2(b)], then in general transitions between the subbands accompanied by optical phonon emission are still possible (as is indicated by the arrow). However, they would not be resonant ones. Indeed, at $eV = \hbar\omega_0$, both the right-hand part of band 2 and the left-hand part of band 1 are occupied so that the resonant transitions are forbidden. In Fig. 2(c), however, the electrostatic potential exceeds the value of that shown in Fig. 2(b), to the extent that $\mu^{(-)} \leq \epsilon_1(0)$. The resonances are allowed but at a voltage bias that is higher than in Fig. 2(a).

A number of cusps appear in Fig. 3(a), each corresponding to a current drop due to electron-phonon interaction. Their sharpness depends on the magnitude of δR, the amount by which EPR condition (1) is mismatched, namely, $\hbar\omega_0 = \epsilon_m(0) - \epsilon_n(0) - \delta R$ (δR can be either positive and negative). In the limit of $\delta R \rightarrow 0$, the sharpness of real resonances is approached. The signature of EPR is particularly pronounced for the differential conductance $g(V)$ and may prove to be instrumental in the experimental investigation of this effect [see Fig. 3(b)].

In the absence of electron-phonon scattering, new uppermost channels are populated with increasing applied voltage bias, each adding a half integer of G_0 to the collisionless differential conductance (see Fig. 1). For such a new uppermost channel m, if the partial current ΔJ_{nm} that flows from channel m to any n is large because it meets EPR condition (1), the total current J no longer increases with increasing voltage bias as $J^{(0)}$ would have, but it decreases instead [see Fig. 3(a)]. Hence resonant drops in current can be attributed to the blockage of the uppermost current carrying channel that would be open in the absence of collisions with phonons at the voltage bias values corresponding to the positions of cusps. The thermal averaging built into the integral Eq. (5) via the Fermi functions and the phonon occupation broadens the range of energies at which the effect of the divergence is evident. The resistance peak associated with EPR is broadened and, consequently, the voltage bias threshold is lowered [17].

As explained above, to the first order in the electron-phonon scattering, the contributions from different phonon branches are additive. The acoustical phonons, having finite occupation numbers at low temperatures, can also be generated in a strong electric field. However, as shown in Ref. [18], because of a larger strength of the coupling with the optical phonons and their larger frequency, one may expect that the main features of the resonant behavior predicted here remain observable.

One can see that actually, exactly as in the case of magnetophonon resonance, the existence of EPR is determined only by the singularities of the electron density of states (DOS) and the existence of a nonvanishing phonon limiting frequency for $q \longrightarrow 0$ and is independent of any other details of the electron and phonon spectrum and their interaction. Thus the existence of the 1D channels, rather than their exact structure, is sufficient for EPR to exist.

In summary, we have calculated a reduction in current due to the electron-phonon scattering in wires connecting two thermal reservoirs with a large voltage bias. We have found oscillations in the resistance as a function of the applied voltage bias. The sharp resonances are a consequence of the confinement and of a special form of the DOS in the 1D case. As electrons undergo the transitions between the bottoms of the subbands, the resonant condition (1) comes into play. We have also pointed out the apparent similarity between the magnetophonon and electrophonon resonances. We have found that EPR is extremely sensitive to the variations of either the gate or bias voltages. Finally, we have found that at finite temperatures the voltage bias threshold required for the onset of EPR may be somewhat lower than the low temperature threshold given by condition (1).

We believe that this effect can be used as a probe for investigating various physical characteristics of nanostructure systems: investigating the character of the actual interactions between the electrons and optical phonons (bulk as well as confined and interface modes), ascertaining the details of the electron band structure and the actual positions of the levels of transverse quantization, and examining the generation of nonequilibrium phonons (as well as the emission of submillimeter electromagnetic waves—cf. with [19]) in the transport phenomena, to name a few. We also encourage experimentalists to think of a physical situation where a blockage of a single channel due to EPR may play a crucial role so that one can separate a one-channel contribution experimentally. Finally, we stress the significance this effect may have on the utilization of nanostructures for the robust applications.

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