Acoustoelectric effect in nanostructures: Role of quasimomentum balance

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In a recent paper giant quantum oscillations of acoustoelectric current under gate voltage variation were predicted for ballistic transport in nanowires. In particular, it was shown that the oscillations should exist provided that the sound frequency ω exceeds a certain threshold value ω_{th} and no acoustoelectric effect exists for $\omega < \omega_{\mathrm{th}}$. The result seems to be in contradiction with a simple energy and quasimomentum balance considerations that in the bulk case lead to the so-called Weinreich relation for the acoustoelectric current. We discuss this paradox in detail and prove it to be related to the very nature of ballistic transport. We develop a theory of the acoustoelectric effect for $\omega < \omega_{\mathrm{th}}$ assuming that the transport is not purely collisionless. [S0163-1829(98)07040-4]

An acoustic wave propagating in a semiconductor creates a net drag of the electrons and hence a dc acoustoelectric current J or, if the circuit is disconnected, a dc acoustoelectric potential difference V. The acoustoelectric effect is a second-order effect in terms of the strength of coupling between the sound and electrons. Recently this effect was considered for the interaction of an ordinary three-dimensional (3D) traveling sound wave with one-dimensional (1D) ballistic electrons.[1] As is well known, the conductance of such an electron system is a steplike function of the Fermi level position[2,3] (which can be monitored by the gate voltage). Each step corresponds to the inclusion of a new mode of transverse quantization to the conduction process and has a height of $G_0 = 2e^2/h$. As for the acoustoelectric current J, in the lowest approximation it has to be expressed in terms of a time average of a bilinear product of the components of elastic displacements or their derivatives. The purpose of the present paper is to investigate theoretically acoustoelectric effect in nanowires where, apart from the electron-sound interaction, there is also a weak scattering of the conduction electrons.

The acoustoelectric current J is a result of the drag of the electrons due to absorption of the ultrasonic phonons and, due to the spatial quantization, it exhibits giant quantum oscillation as a function of the Fermi level position that can be monitored by the gate voltage. This phenomenon is similar to giant quantum oscillations of the ultrasonic absorption in a magnetic field[4] in bulk samples where due to a system of Landau levels analogous oscillations can be observed in the course of magnetic field variation. An important feature of the predicted effect in nanostructures is vanishing of the acoustoelectric current for acoustic frequencies below the threshold value[1]

$$\omega_{\mathrm{th}} = 2mw^2/\hbar. \tag{1}$$

Here m and w are the electron effective mass and sound velocity, respectively.

Keeping in mind that the corresponding acoustic absorption does not vanish for $\omega < \omega_{\mathrm{th}}$, one sees a contradiction with the physical picture known for the acoustoelectric effect in the bulk structures. For the latter the acoustoelectric current density j is known to obey the well-known Weinreich relation:[5]

$$j = \mathcal{M}\frac{\Gamma S}{w}. \tag{2}$$

Here $\mathcal{M}$ is the electron mobility, Γ is the acoustic absorption coefficient ($1/\Gamma$ has units of length), and S is the acoustic intensity. As is easily seen, ΓS is the total rate of absorption of the acoustic energy density by electrons and thus $\Gamma S/w$ is the rate of quasimomentum transfer from the acoustic wave to the electrons per one second per unit volume. Thus the drag (acoustoelectric) current is directly related to the quasimomentum transfer and, in general, the acoustoelectric current is expected to exist whenever the ultrasound is absorbed by the conduction electrons and therefore the quasimomentum transfer is nonzero.

In contrast, for the ballistic case considered in Ref. 1, for the sound of comparatively low frequencies, the drag current vanishes even if the acoustic absorption, and therefore the quasimomentum transfer from the acoustic flux to electrons, is nonzero. At first the result seems to be paradoxical. In what follows we show that the paradox is a direct consequence of the ballistic nature of the transport. We will also show that an inclusion of a weak additional (''background'') scattering leads to a nonvanishing drag current. However, if the scattering is weak the Weinreich relation is not restored. One gets instead a new relation between the acoustic flux and acoustoelectric current density [see below Eqs. (27), (28), and (35)].

Following the approach used in Ref. 1, we begin by giving a brief review of the main results exhibiting the paradox. We represent the ultrasound as a flux of phonons with the same frequency $\omega_{\mathbf{q}}$ and wave vector $\mathbf{q}$. The phonons are assumed to propagate in the bulk of the sample, i.e., to be three dimensional (3D). The nanostructure in the form of a quantum wire of a constant cross section is oriented along the x axis. This means that in the course of electron-phonon collisions the x component of the quasimomentum is conserved whereas q_y and q_z are not because the electrons are confined in y and z directions.

We will assume that the direction of phonon propagation, i.e., their wave vector $\mathbf{q}$, is also parallel to the x axis. Then the energy and quasimomentum conservation law takes the form

$$\epsilon_n(p)+\hbar\omega_{\mathbf{q}}=\epsilon_n(p+\hbar q). \tag{3}$$

Here $\epsilon_n(p)=\epsilon_n(0)+p^2/2m$ is an energy of electron belonging to a one-dimensional (1D) subband (channel), n is the quantum number of transverse quantization, $\epsilon_n(0)$ is the position of the nth level of transverse quantization, p is the x component of electron quasimomentum, while the phonon frequency is $\omega_{\mathbf{q}}=wq$ where w is the sound velocity. $p=mw-\hbar q/2$, and $p+\hbar q=mw+\hbar q/2$ are the solutions of Eq. (3) where we have omitted the index x for the phonon wave vector q. The quasimomentum transfer from phonons to electrons, $\hbar\mathbf{q}$, is expected to bring about an acoustoelectric current J across the nanowire.

The initial and final quasimomenta of electrons taking part in such transitions are completely determined by Eq. (3). Such transitions occur if the initial electron state is either within a thermal layer near the Fermi level (if $\hbar\omega_{\mathbf{q}}\ll k_{\mathrm{B}}T$) or within a layer of width $\hbar\omega_{\mathbf{q}}$ between $\mu_n-\hbar\omega_{\mathbf{q}}$ and μ_n (assuming that $\hbar\omega_{\mathbf{q}}\gg k_{\mathrm{B}}T$), where $\mu_n=\mu-\epsilon_n(0)$, μ being the chemical potential.

Consider, for instance, the case $\hbar\omega_{\mathbf{q}}\gg k_{\mathrm{B}}T$. When in the course of gate voltage variation both the initial and final states appear occupied or unoccupied the acoustoelectric current drops. With further change in the gate voltage an electron with its initial state belonging to another subband n' moves into the layer between $\mu_{n'}-\hbar\omega_{\mathbf{q}}$ and $\mu_{n'}$, which leads (provided that its final state is outside the layer, i.e., is unoccupied) to an increase in acoustoelectric current. This results in the giant quantum oscillation of acoustoelectric current as a function of the gate voltage.

In the absence of sound, the distribution function of electrons, f_{pn}, is simply the Fermi function, $f^{(\mathrm{F})}(\epsilon)$, where $\epsilon=\epsilon_n(p)$ is the electron energy. Allowing for a weak interaction between electrons and the ultrasonic phonons results in $f_n=f^{(\mathrm{F})}(\epsilon_n)+\Delta f_n$ with Δf satisfying the equation

$$v\frac{\partial\Delta f}{\partial x}=I[f]. \tag{4}$$

Here $v=\partial\epsilon_n/\partial p$ is the electron velocity (which does not depend explicitly on the quantum number n). $I[f]$ is the electron-phonon collision term. We have dropped from the right-hand side of Eq. (4) the term $e(\partial\Delta\phi/\partial x)(\partial f^{(F)}/\partial p)$ (where $\Delta\phi$ is the time averaged electrostatic potential) as it gives no contribution to the current. We will drop the 1D subband (channel) index n wherever it will not create confusion.

We start with the analysis of the simplest boundary condition $\Delta f=0$, which is satisfied at $x=\mp L/2$ for $p>0$ ($p<0$), respectively. (Below we will consider a more general boundary condition.) For $p>0$ ($p<0$), respectively, the solution of Eq. (4) is

$$\Delta f(x)=(x\pm L/2)\frac{1}{v}I[f]. \tag{5}$$

Here $x=0$ lies at the midpoint of a conductor of a total length L. $\Delta f(x)$ is proportional to $1/v$. This is physically natural as for the ballistic transport the nonequilibrium part of the distribution function relaxes because the conduction electrons leave the quantum wire in the course of their motion. The smaller the x componenet of the electron velocity, the slower the relaxation. The explicit form of the collision term reads (we assume that only transitions within one 1D electron band are allowed by the energy and quasimomentum conservation)

$$I[f]=\int\frac{dp'}{2\pi\hbar}\int\frac{d^2q_\perp}{(2\pi)^2}W(f'-f)N$$
$$\times[\delta(\varepsilon'-\varepsilon-\hbar\omega_{\mathbf{q}})+\delta(\varepsilon'-\varepsilon+\hbar\omega_{\mathbf{q}})], \tag{6}$$

where N is the phonon distribution function (we have discarded 1 since $N\gg1$), W is the coefficient of electron-phonon interaction introduced in Ref. 1. For the deformation potential interaction we have

$$W_{\mathbf{q}}=\pi\Lambda^2q^2/\rho\omega_{\mathbf{q}}, \tag{7}$$

where Λ is the deformation potential constant for the phonon branch in consideration and ρ is the mass density. For the unscreened piezoelectric interaction

$$W_{\mathbf{q}a}=\frac{\pi}{\rho\omega_{\mathbf{q}}}\left[\frac{4\pi e\beta_{q,lq}\nu_l(\mathbf{q},a)}{\varepsilon_{qq}}\right]^2. \tag{8}$$

Here $\beta_{i,ln}$ is the tensor of piezoelectric moduli (which is symmetric in the last two indices), ε_{il} is the tensor of dielectric susceptibility, $\nu_l(\mathbf{q},a)$ is the polarization vector (i.e., a unit vector along the elastic displacement $\mathbf{u}$) of the phonon with wave vector $\mathbf{q}$, belonging to the branch a. Index q indicates the projection of a tensor on the $\mathbf{q}$ direction.

The total current is given by

$$J=2e\sum_n\int_{-\infty}^{\infty}\frac{dp}{2\pi\hbar}v\Delta f_{np}(x), \tag{9}$$

where $\Delta f_{np}(x)$ is given by Eq. (5). The integrand in Eq. (9) is proportional to electron group velocity $v=p/m$ whereas $\Delta f_{np}(x)$ according to Eq. (5) is proportional to v^{-1}. As a result we have

$$J=2eL\sum_n\int_0^{\infty}\frac{dp}{2\pi\hbar}I[f]-2eL\sum_n\int_{-\infty}^{0}\frac{dp}{2\pi\hbar}I[f]$$
$$+4ex\sum_n\int_{-\infty}^{\infty}\frac{dp}{2\pi\hbar}I[f]. \tag{10}$$

The last term in Eq. (10) vanishes since collisions do not change the total number of electrons.

The integrations in Eq. (6) are in fact over the three components of the phonon wave vector. The 3D phonon distribution function $N_{\mathbf{k}}$ can be presented in the form

$$N_{\mathbf{k}}=[(2\pi)^3 S/\hbar w^2 k]\delta^{(3)}(\mathbf{k}-\mathbf{q}), \tag{11}$$

where S is the sound intensity. For the electron distribution function Δf we have Eq. (5) with

$$I=A[\delta(p-p_2)-\delta(p-p_1)]. \tag{12}$$

Here

$$p_{1,2}=mw\mp\hbar q/2 \tag{13}$$

and

$$A=\frac{SWm}{(\hbar\omega_{\mathbf{q}})^2}(f^{(\mathrm{F})}_{mw-\hbar q/2}-f^{(\mathrm{F})}_{mw+\hbar q/2}). \tag{14}$$

Inserting Eq. (11) into Eq. (6) and making use of Eq. (10) one gets for $\hbar q>2mw$

$$J=\frac{emSWL}{2\pi\hbar^3\omega_{\mathbf{q}}^2}\sum_n [f^{(F)}(\epsilon^{(-)}-\mu_n)-f^{(F)}(\epsilon^{(+)}-\mu_n)], \tag{15}$$

where $\epsilon^{(\pm)}=(\hbar q/2\pm mw)^2/2m$, and we remind that $\mu_n=\mu-\epsilon_n(0)$.

Taking into account the δ-function structure of the collision term [Eq. (12)] one can check explicitly the fact that the electron-phonon collisions conserve the number of electrons

$$\left[\frac{\partial n}{\partial t}\right]_{\mathrm{coll}}=2\sum_n\int_{-\infty}^{\infty}\frac{dp}{2\pi\hbar}I[f] =\sum_n\int_0^{\infty}\frac{dp}{\pi\hbar}I[f]+\sum_n\int_{-\infty}^{0}\frac{dp}{\pi\hbar}I[f]=0. \tag{16}$$

Since the integrand of Eq. (10) is proportional to the sum of delta functions one can write

$$\sum_n\left(\int_0^{\infty}\frac{dp}{2\pi\hbar}I[f]-\int_{-\infty}^{0}\frac{dp}{2\pi}I[f]\right) =\frac{1}{2}2\frac{1}{\hbar q}\sum_n\int_{-\infty}^{\infty}\frac{dpp}{2\pi\hbar}I[f]=\frac{1}{2\hbar q}\frac{\partial\mathcal{P}}{\partial t}$$

where $\partial\mathcal{P}/\partial t$ is the overall rate of quasimomentum transfer from the phonon flux to electrons. Correspondingly, one gets the following relation between the acoustoelectric current and the rate of quasimomentum transfer to the electrons:

$$J=\frac{eL}{\hbar q}\frac{\partial\mathcal{P}}{\partial t}. \tag{17}$$

This is a sort of Weinreich relation for 1D electrons (and 3D phonons) while the factor $eL/\hbar q$ plays a role of the ''effective mobility.'' As $\hbar q/m$ is of the order of the velocity of electrons interacting with the acoustic wave the ''effective relaxation time'' $mL/\hbar q$ appears to be of the order of the time of flight of an electron across the nanostructure, as may be indeed expected for the ballistic structures.

However, for $\hbar q<2mw$ (or $\omega<\omega_{\mathrm{th}}$) the situation differs drastically. Indeed, when $\hbar q<2mw$ and therefore both $mw-\hbar q/2$ and $mw+\hbar q/2$ are positive the δ functions do not contribute to the integral

$$\int_{-\infty}^{0}\frac{dp}{\pi\hbar}I[f]$$

in Eq. (16). So we are left with

$$\left[\frac{\partial n}{\partial t}\right]_{\mathrm{coll}}=\frac{J}{eL}=2\sum_n\int_0^{\infty}\frac{dp}{2\pi\hbar}I[f]. \tag{18}$$

In view that the integral (18) should vanish because the collisions conserve the concentration of electrons, one sees that the acoustoelectric current, which is proportional to the same integral, should also vanish.

Yet the rate of quasimomentum transfer in this case is given by

$$\frac{\delta\mathcal{P}}{\delta t}=2\sum_n\int_0^{\infty}\frac{dp}{2\pi\hbar}pI[f]=\frac{Aq}{\pi}. \tag{19}$$

One sees that this quantity (directly related to the number of acoustic quanta absorbed by electrons) does not vanish, unlike the current. Thus we arrive at a contradiction with the physics of quasimomentum balance discussed above. The resolution of the contradiction is related to the very nature of the ballistic transport.

On the one hand, for $\omega<\omega_{\mathrm{th}}$, the electron-phonon collisions do not change the direction of electron quasimomentum. Thus only the electrons penetrating into the wire from one of the reservoirs (say, the ''left'' one) are coupled to the phonons while those penetrating from the ''right'' one are insensitive to the sound. Therefore we have the boundary problem formulated only for the ''left'' boundary at $x=-L/2$. On the other hand, the total number of the electrons passing through this boundary is evidently not affected by the acoustic wave. Therefore, the current of the ''left'' electrons is completely controlled by the corresponding boundary condition keeping it equal to a countercurrent of the ''right'' electrons.

This situation can be described in different words. The transitions due to the absorptions of acoustic quanta are allowed only between the electrons with positive velocities $(v>0)$. The contribution of each such electron to the current is, on the one hand, *proportional* to v. On the other hand, it is proportional to the time it spends within the quantum wire, which is *inversely proportional* to v. As a result, it is v independent and the absorption of ultrasonic quanta does not change the current although it enhances the total electron quasimomentum $\mathcal{P}$.

While the total current vanishes, the acoustic wave does affect the local concentration of the electrons. Indeed, as can be obtained from Eq. (5):

$$\Delta n = \left(x + \frac{L}{2}\right) \frac{Sm^2 W}{\hbar q (\hbar \omega_{\mathbf{q}})^2} \left[\left(\frac{mw}{\hbar q}\right)^2 - \frac{1}{4}\right]^{-1} \times \sum_l \left[f^{(F)}(\epsilon^{(-)} - \mu_l) - f^{(F)}(\epsilon^{(+)} - \mu_l)\right]. \quad (20)$$

Here we neglect the screening effects. If the latter were strong and supported the neutrality of the system, one would expect a formation of the screening electrostatic field which, in its turn, would create a nonzero acoustoelectric current.

If, in addition to the interaction with the acoustic wave, one takes into account other scattering processes where the conduction electrons take part one will get a nonzero acoustoelectric current. To begin with, we will consider the simplest form of elastic processes, i.e. the reflection of electrons at both ends of the quantum wire (cf. with Ref. 6 where scattering by a long-range potential due to the edges of the channel has been considered). The reflection coefficients will be assumed the same at both ends of the wire and will be denoted R_1 and R_2 for p_1 and p_2, respectively. The role of such processes is particularly interesting for $\hbar q/2mw < 1$ where without these processes one gets $J=0$. Let us turn to this case. In other words, let us assume that p_1 and p_2 are positive.

The transport equation (4) can be now rewritten [with account of Eq. (12)] as

$$v \frac{\partial \Delta f_p}{\partial x} = A[\delta(p-p_2) - \delta(p-p_1)]. \quad (21)$$

The boundary condition at $x=-L/2$ now has the form

$$\Delta f_p = R \Delta f_{-p}, \quad (22)$$

so that the solution of Eq. (21) is

$$\Delta f_p = R(p) \Delta f_{-p} + \left(x + \frac{L}{2}\right) \frac{m}{p} A[\delta(p-p_2) - \delta(p-p_1)]. \quad (23)$$

Below we will assume that $p>0$, so that the distribution function for the negative values of p we will denote by $f(-p)$.

Now, the boundary condition at $x=L/2$ has the form

$$\Delta f_{-p} = R \Delta f_p. \quad (24)$$

Making use of Eq. (23) one can write Eq. (24) as

$$\Delta f_{-p} = R^2 \Delta f_{-p} - RL \frac{m}{p} A[\delta(p-p_2) - \delta(p-p_1)], \quad (25)$$

so that

$$\Delta f_{-p} = -L \frac{m}{p} A\left[\frac{R_2}{1-R_2^2} \delta(p+p_2) - \frac{R_1}{1-R_1^2} \delta(p+p_1)\right], \quad (26)$$

where $R_{1,2} = R(p_{1,2})$. We remind that Δf_{-p} is not perturbed directly by the acoustic wave.

As is mentioned above, positive values of p give no contribution to the acoustoelectric current. The contribution of $p<0$ is given by

$$J = e \int_{-\infty}^{0} \frac{2dp}{2\pi\hbar} \Delta f_{-p}$$

(we will assume that only one 1D electron band can contribute to the acoustoelectric current and will drop summation over n in such expressions). Making use of Eq. (26) we get

$$J = L \frac{eSWm}{\pi \hbar^3 \omega_q^2} \left(f^{(\mathrm{F})}_{mw-\hbar q/2} - f^{(\mathrm{F})}_{mw+\hbar q/2}\right) \left(\frac{R_1}{1-R_1^2} - \frac{R_2}{1-R_2^2}\right). \quad (27)$$

This equation represents giant quantum oscillations of the acoustoelectric current whose amplitude is smaller than in the collisionless case [Eq. (15)] by the factor

$$\frac{R_1}{1-R_1^2} - \frac{R_2}{1-R_2^2}.$$

It is also interesting to consider an asymmetric case where the reflection coefficient for the electrons with $p<0$ is zero at $x=-L/2$ and is nonvanishing and p dependent only at $x=L/2$. Then one has

$$J = L \frac{eSWm}{\pi \hbar^3 \omega_q^2} \left(f^{(\mathrm{F})}_{mw-\hbar q/2} - f^{(\mathrm{F})}_{mw+\hbar q/2}\right)(R_1 - R_2), \quad (28)$$

where $R_{1,2}$ are the values of the electron reflection coefficients for $x=L/2$ and quasimomenta p_1 and p_2, respectively. In the same manner one can consider a more general case where the scatterer has some intermediate coordinate x_0 within the quantum wire.

Thus we see that even additional scattering of a simplest form results in an acoustoelectric current though the Weinreich relation in its usual form does not hold for this case. We also have physical considerations that in a nonlinear case where the sound intensity is so high that there is no direct proportionality between the intensity and acoustoelectric current one can expect the existence of acoustoelectric current even for $\hbar q < 2mw$. The nonlinear case, however, deserves a special treatment (cf. with Refs. 7 and 8).

Let us now consider another limiting case where the number of scatterers is large enough while their spatial distribution along the channel is uniform, which corresponds to the standard expression for the collision integral. The transport equation (4) can be now rewritten [with account of Eq. (12)] as

$$v \frac{\partial \Delta f}{\partial x} + \Delta f(p) \int \frac{dp'}{2\pi\hbar} P_{pp'} - \int \frac{dp'}{2\pi\hbar} P_{p'p} \Delta f(p') = A[\delta(p-p_2) - \delta(p-p_1)]. \quad (29)$$

Here $P_{p'p}$ is the probability of elastic scattering from state p to state p', while the quantities $p_{1,2}$ are introduced above. Here we assume the scatterers to be short-range ones. We would like also to note that for the purely 1D case the probabilities $P_{p'p}$ actually describe the backscattering with $p'=-p$ and thus $P_{p'p}\equiv P(|p|)\delta(p'+p)$.

The second term on the left-hand side of Eq. (29) is the so-called ''out'' term of the collision operator that can be presented as $\Delta f_p/\tau_p$, where

$$\frac{1}{\tau_p}\equiv\frac{P(|p|)}{2\pi\hbar} \tag{30}$$

is the electron relaxation time due to elastic collisions. The third term on the left-hand side of Eq. (29) is the ''in'' term.

The boundary conditions will be assumed to have the simplest form, namely,

$$\Delta f|_{x=\mp L/2}=0$$

for $p>0$ and $p<0$, respectively.

Assuming the elastic scattering to be weak we will use iterations to solve Eq. (29). In the first approximation one neglects the elastic scattering. For the case $p_1>0$ and $p_2>0$ one has, for the region $-L/2<x<L/2$ with regards to the ''zero'' boundary condition at $x=-L/2$,

$$\Delta f_1=\left(x+\frac{L}{2}\right)A\left[\frac{1}{v_2}\delta(p-p_2)-\frac{1}{v_1}\delta(p-p_1)\right]. \tag{31}$$

Now we make the next iteration in the scattering probabilities $P_{p'p}$. We will have contributions proportional to the ''out'' and ''in'' terms, respectively. The ''out'' contribution naturally has the same structure as the perturbation and is quadratic in x:

$$\Delta f_2^{(\mathrm{out})}(p)=-\frac{1}{2}\left(x+\frac{L}{2}\right)^2\frac{A}{\tau_{\mathbf{p}}}\left[\frac{1}{v_2^2}\delta(p-p_2)-\frac{1}{v_1^2}\delta(p-p_1)\right]. \tag{32}$$

The ''in'' contribution is related only to negative values of v (because one deals with a backscattering of electrons with positive values of the initial velocity). Thus the boundary condition should be formulated for $x=L/2$ and one finally obtains

$$\Delta f_2^{(\mathrm{in})}(-p)=\frac{A}{2\tau}\left[L^2-\left(x+\frac{L}{2}\right)^2\right][\delta(p-p_2)-\delta(p-p_1)]. \tag{33}$$

To calculate the current related to Δf_2 one should multiply this quantity by v and sum over p. For the contribution of the ''in'' term one can change the variable in this integration $p\to -p$, and the term $\propto(x+L/2)^2$ is seen to cancel out the contribution of the ''out'' term.

Thus

$$J=\frac{eLSWm}{2(\hbar\omega_{\mathbf{q}})^2}(f^{(\mathrm{F})}_{mw-\hbar q/2}-f^{(\mathrm{F})}_{mw+\hbar q/2})\int_0^\infty\frac{v}{\tau_p}\times\left[\frac{1}{v_1^2}\delta(p-p_1)-\frac{1}{v_2^2}\delta(p-p_2)\right]\frac{2dp}{2\pi\hbar}. \tag{34}$$

We write Eq. (34) in this form to emphasize that in this case the acoustoelectric current is proportional to the relaxation rate rather than to the relaxation time as is the case in Ref. 5. Physically this means that here [as above—see Eq. (28)] the acoustoelectric current is determined by a single elastic scattering event.

Making use of the δ functions in the integrand one can present Eq. (34) in the following final form:

$$J=\frac{eSWmL}{2\pi\hbar^3\omega_{\mathbf{q}}{}^2}(f^{(\mathrm{F})}_{mw-\hbar q/2}-f^{(\mathrm{F})}_{mw+\hbar q/2})L\left(\frac{1}{\tau_1 v_1}-\frac{1}{\tau_2 v_2}\right), \tag{35}$$

where $\tau_{1,2}=\tau(p_{1,2})$. The products $L/\tau_i v_i$ should be considered as small parameters of our theory.

It is interesting to compare our results with those obtained by Shilton *et al.*[9] where the acoustoelectric current in the 1D channel (and the corresponding giant quantum oscillations effect) was examined. For the nonballistic case $v\tau_i<L$ (where τ_i is the relaxation time with respect to the impurity scattering introduced in Ref. 9) the final result is proportional to τ_i (i.e., to the mobility $\mathcal{M}$) and does not vanish for any acoustic frequency. Such behavior is typical for a situation described by the Weinreich relation.[5] However, for the ballistic or ''almost'' ballistic situation $v\tau_i>L$, a nonvanishing current is also predicted in Ref. 9 for all the acoustic frequencies, which is in disagreement with our results. In our opinion, this difference is related to the fact that in Ref. 9 the inverse relaxation operator approximation was extended to the ballistic situation (namely, a concept of ''escape rate'' $\sim v/L$ was used). We believe that the very nature of the ballistic transport may be not always compatible with the use of this approximation. We stress that the pure ballistic transport is related to a boundary condition problem that, as was shown, for the low acoustic frequencies results in *vanishing of the current*.

In summary, we have analyzed acoustoelectric current in a nanowire under the quasiballistic conductance regime. While for comparatively large frequencies in the absence of electron scattering the current obeys a sort of relation determined by the quasimomentum conservation, for comparatively low sound frequencies the current vanishes (in the absence of screening) despite the nonzero quasimomentum transfer from the acoustic flux to the electrons. This paradox is shown to be related to the very nature of the ballistic transport where for small frequencies the sound wave appears to be coupled only to electrons entering the structure from one of the two reservoirs connected by the nanostructure, the total electron flux being fixed by the boundary condition. We have shown that the effect of weak scattering leads to the restoration of the acoustoelectric current for the low-frequency limit. An explicit equation for the acoustoelectric current is obtained for this case. In this case too we predict giant quantum oscillations of the acoustoelectric current.

The existence of such an effect is, in essence, based on the conservation laws for the energy and x component of the quasimomentum in the electron-phonon interaction. Therefore it should exist not only for the 3D bulk sound but also for the 2D surface acoustic waves interacting with a quantum wire near the surface of the sample (cf. with Ref. 9). The effect can be used for detection of the ultrasound, for investigation of the electron spectrum in microstructures, and (which could prove to be one of the most interesting applications) for investigation of ''background'' scattering of the electrons in nanostructures.

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