Photovoltage in curved one-dimensional systems

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Curvature of quantum wire results in the intrasubband absorption of IR radiation that induces stationary photovoltage in the presence of circular polarization. This effect is studied in ballistic (collisionless) and kinetic regimes. The consideration is concentrated on quantum wires with a curved central part. It is shown that, if mean-free path is shorter than length of the curved part, the photovoltage does not depend on the wire shape, but on the total angle of rotation of the tangent to the wire. It is not the case when the mean-free path is finite or large. This situation was studied for three specific shapes of wires: "hard angle," "open book," and "Ω-like."

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I. INTRODUCTION

The stationary current induced by alternating force was a subject of numerous publications. The ordered motion of electrons in the absence of stationary driving force implies the simultaneous existence of energy and momentum sources. While the energy originates from the alternating force, the momentum can be transferred from the wave itself, as well as (constructively) from scattering events, when vector asymmetry dictates the direction of the current.

The directions of study can be classified in relation to the participation of light momentum, the source of anisotropy, the system uniformity, coherence of light, etc. In particular, the term "photogalvanic effect"[1–3] is used to describe the stationary photocurrent in a homogeneous medium with low symmetry, where the direction of current is determined by the electric field of light together with the third-rank tensor belonging to the medium itself, while the directional motion is caused by the participation of electron scattering. The term "photon drag"[4–6] relates to currents due to the transmission of momentum from photons to electrons and does not need any participation of the "third body," namely, scatterers. The term "quantum pumps"[7–10] is applied mainly to the local quantum systems driven by periodically changing parameters. The term "ratchet"[11–13] is used to describe the stationary flow caused by an alternating force, not necessarily electrical.

The purpose of the present paper is IR photoresponses in curved quantum wires. The main idea is that, in spite of uniformity of IR electric field, the acting component of electric field tangential to the narrow wire becomes nonuniform with the characteristic length dictated by the wire curvature. This idea reminds of the situation with the nonuniformity of acting component of the magnetic field in the curved two-dimensional (2D) systems (see, e.g., Ref. 14). We studied the photocurrent in a spiral quantum wire earlier.[15] The system under examination differs from Ref. 15 in the nonhomogeneous curvature. An example of the nonuniformly curved system, namely, curved one-dimensional (1D) quantum dot lattice subjected to IR radiation was studied in Ref. 16.

We consider planar quantum wires with kinks subjected to normally incident arbitrarily polarized electromagnetic wave. A sketch of the considered system and examples of curved wires are shown in Figs. 1 and 2. The wire is assumed to be strictly one dimensional when only the tangential component of the electric field (inhomogeneous due to curvature) affects electrons in a classical manner. This produces a stationary current in closed circuit (shown in Fig. 1) or a voltage in open-circuit regimes. The problem is studied in the framework of the classical Boltzmann kinetic equation for freely moving electrons along the curved quantum wire (the case of a single subband occupied with intrasubband absorption). We assume that the lowest subband bottom of the wire is flat. It means we ignore the inhomogeneities caused by the bottom potential, wire width, or curvature.

In the lowest order on intensity the stationary current is the second-order response to electric field $\mathbf{E}(t)$. In particular, for coherent monochromatic radiation $\mathbf{E}(t)=\mathrm{Re}(\mathbf{E}^{\omega}e^{-i\omega t})$, $\mathbf{E}^{\omega}=E^{\omega}(1,i\zeta)/\sqrt{2}$ is the complex amplitude of wave; the degree of circular polarization $\zeta=0$ for linear polarized and $\zeta=\pm 1$ for fully circularly polarized wave. Generally speaking, the global current can be connected with the electric field amplitude in the lower order by a tensor relation $J_i=\mathrm{Re}(\alpha_{ijk}E_j^{\omega}E_k^{-\omega})$, where α_{ijk} is some tensor determined by the shape of the curve.[17] The plane curved wires considered here, symmetric with respect to $x\leftrightarrow -x$, can be globally characterized by a unit vector $\mathbf{b}\parallel Oy$. (See Fig. 1 where up and down directions are not equivalent.) In that case the tensor α_{ijk} (and the direction of the stationary current $\mathbf{J}$) can be constructed by means of vector $\mathbf{b}$ as follows:

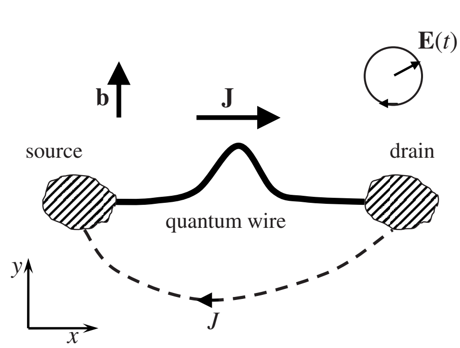

FIG. 1. Sketch of the considered system. Quantum wire curved in its central part connects source and drain. The external alternating electric field $\mathbf{E}(t)$ is circularly polarized in the wire plane (x,y).

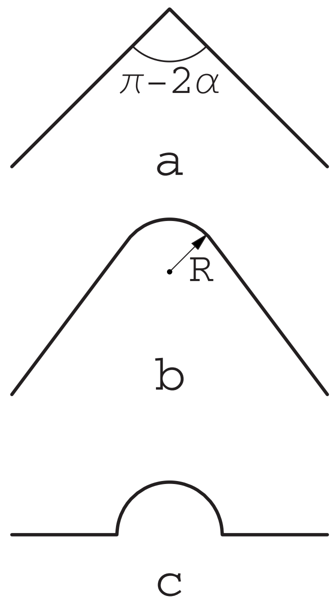

FIG. 2. Considered quantum wires: (a) angle, (b) open book (straight lines tangential to a circular segment), and (c) Ω curve (semicircle with straight source and drain). In cases (a) and (b) the angle between straight parts of wire equals $\pi-2\alpha$ and the radius of circular segments in cases (b) and (c) equals R.

$$\alpha_{ijk} = i\alpha_a(\delta_{ij}b_k - \delta_{ik}b_j) + \alpha_s[(\delta_{ij}b_k + \delta_{ik}b_j)/2 - b_ib_jb_k]$$

and $\mathbf{J}=i\alpha_a[\mathbf{b}[\mathbf{E}^\omega\mathbf{E}^{-\omega}]]+\mathrm{Re}[\alpha_s(\mathbf{E}^\omega-\mathbf{b}(\mathbf{E}^\omega\mathbf{b}))(\mathbf{b}\mathbf{E}^{-\omega})]$, where α_a and α_s are real parameters. As will be seen below, only the term with α_a in the current survives in the considered approximation. Since $[\mathbf{E}^\omega\mathbf{E}^{-\omega}]=i\zeta(E^\omega)^2$, the current should change sign with the sign of circular polarization and should vanish for linear polarization. So this effect can be attributed to the class of circular photogalvanic effects.

Generally, the value of current is determined by the wire shape and its symmetry. Unlike photogalvanic effects (appearing in homogeneous media) the considered effect has local nature, which makes it related to quantum pumps. The current is caused by the momentum transfer from the wire to electrons via inhomogeneity of acting alternating force. From this point of view the considered phenomenon is some sort of the photon drag effect.

Both homogeneity of the system and smallness of electromagnetic field wave vector, especially in the case of far IR light, restrict the known photocurrents. Smallness of the wave vector leads to relatively weak photon drag effect; the photogalvanic effect caused by participation of the scattering by a third body (impurities and phonons) is also weak. Curved 1D or 2D system present a possibility to overcome this weakness: the uniform external field affects propagating electrons by the nonuniform effective force whose characteristic length, in principle, is comparable with the Fermi wavelength.

II. BASIC EQUATIONS

Let us consider the one-dimensional quantum wire of length L described by a planar curve $\mathbf{a}(s)=(a_x,a_y)$, where s is the arc distance along the curve, $-L/2<s<L/2$. The effective electric field $\mathcal{E}$ affecting electrons is determined by the projection of electric field $\mathbf{E}(s,t)$ to unit tangent vector $\mathbf{t}(s)=\mathbf{a}'(s)$: $\mathcal{E}(s,t)=\mathbf{E}(t)\mathbf{t}(s)$. We will consider the effect in the approximation of the classical kinetic equation

$$\frac{\partial f}{\partial t} + v\frac{\partial f}{\partial s} - e\mathcal{E}\frac{\partial f}{\partial p} = -\frac{f-\langle f\rangle}{\tau}, \tag{1}$$

where $\tau=1/\nu$ is the relaxation time and $\langle f\rangle=[f(p)-f(-p)]/2$. We neglect the static potential in the wire, in particular, connected with the curvature itself. Further we consider the effect of the monochromatic electric field $\mathcal{E}(t)=\mathrm{Re}(\mathcal{E}^\omega e^{-i\omega t})$.

We will deal with the specific cases of symmetric planar curves depicted in Fig. 2. The corresponding unit tangent vectors are

$$t_x(s) = \cos\alpha,$$

$$t_y(s) = -\,\mathrm{sign}(s)\sin\alpha \quad (\text{``angle''}), \tag{2}$$

$$t_x(s) = [\theta[-(s+R\alpha)] + \theta(s-R\alpha)]\cos\alpha + \theta(R\alpha - s)\theta(R\alpha+s)\cos\left(\frac{s}{R}\right),$$

$$t_y(s) = -\,\mathrm{sign}(s)\{\theta[-(s+R\alpha)] + \theta(s-R\alpha)\}\sin\alpha - \theta(R\alpha-s)\theta(R\alpha+s)\sin\left(\frac{s}{R}\right) \quad (\text{open book}), \tag{3}$$

$$t_x(s) = \theta\left(-s-\frac{\pi}{2}R\right) + \theta\left(s-\frac{\pi}{2}R\right) + \theta\left(\frac{\pi}{2}R-s\right)\theta\left(\frac{\pi}{2}R+s\right)\cos\left(\frac{s}{R}\right),$$

$$t_y(s) = \theta\left(\frac{\pi}{2}R-s\right)\theta\left(\frac{\pi}{2}R+s\right)\sin\left(\frac{s}{R}\right) \quad (\text{``}\Omega\ \text{curve''}). \tag{4}$$

The coordinate s is counted from the center of the curves.

Despite simple problem formulation, the target setting has pitfalls. In fact, the absence of effective driving force apart from the curved part of the wire results in vanishing of stationary current in this section of a wire and, consequently, due to continuity equation, the current in any part of the wire. One can make certain of that fact directly from Eq. (1) by integration with respect to the momentum.

We shall consider separately two cases: ballistic and kinetic ones. By a ballistic wire we mean a wire shorter than the mean-free path, ends of which join the electron seas. In that case one can neglect the scattering inside the wire and consider entering electrons as equilibrium.

On the contrary, if mean-free path is small as compared with the wire length, the equilibrium is achieved inside the wire while equilibrium states can differ in the right and the left sides of the wire. This case needs to be solved by means of the kinetic equation accounting for the collisions.

III. BALLISTIC WIRE

In the ballistic limit we start from Eq. (1) with omitted right-hand side. Solving the kinetic equation in the second order of electromagnetic field we find the correction to the isotropic part of the distribution function. From the kinetic equation it follows that the isotropic part of the distribution function should have different values on the ends of the wire $s=\pm L/2$. The difference of two limits $[\int dp\, f(p)_{s=L/2} - \int dp\, f(p)_{s=-L/2}]/\pi$ can be attributed to the difference of concentrations on the contacts Δn, in other words, to the difference of their chemical potentials $\Delta\mu=\Delta n/(\partial n/\partial\mu)$. Taking into account the electroneutrality reasons we should keep the concentration on the wire ends that can be done by adding a static voltage such as $eV=\Delta\mu$.

The solution of linearized collisionless kinetic equation satisfying the condition of equilibrium on the wire ends reads as

$$f_1^\omega = \exp\left(\frac{i\omega s}{v}\right)\int_{\mp L/2}^{s} ds' \exp\left(\frac{-i\omega s'}{v}\right) e\mathcal{E}^\omega(s') f_0', \quad (5)$$

where f_0 is the Fermi function, prime means the derivation over energy $\varepsilon=p^2/2m$, and L is the normalization length. The upper (lower) sign in the limit of the integral corresponds to $v>(<)0$.

For Δn we have

$$\Delta n = \frac{1}{\pi}\int_{-\infty}^{\infty} dp\, \bar{f}_2 = \frac{e}{2\pi m}\mathrm{Re}\int_{-\infty}^{\infty}\frac{dp}{v^2}\int_{-\infty}^{\infty} ds\, \mathcal{E}^{-\omega}(s) f_1^\omega(s), \quad (6)$$

where $\bar{f}_2$ is the stationary part of the quadratic in $\mathcal{E}$ distribution function. Inserting expression (5) into Eq. (6) one can obtain

$$\Delta n = -\frac{e^2}{2\pi m}\int_0^\infty \frac{dp}{v^2} f_0' \int_{-L/2}^{L/2}\int_{-L/2}^{L/2} ds ds' \times \sin\left(\frac{i\omega(s-s')}{v}\right)\mathrm{Im}(\mathcal{E}^{-\omega}(s)\mathcal{E}^\omega(s')). \quad (7)$$

Taking into account the symmetry of curves under consideration we obtain for voltage V from Eq. (7)

$$V = V_0\zeta\frac{\omega^2}{\pi(\partial n/\partial\mu)}\int_0^\infty \frac{dp}{v^2} f_0' \int_{-L/2}^{L/2}\int_{-L/2}^{L/2} ds ds' \times \sin\left(\frac{i\omega(s-s')}{v}\right) t_x(s) t_y(s'). \quad (8)$$

Here

$$V_0 = \frac{e}{2m\omega^2}(E^\omega)^2.$$

Introducing the space Fourier transforms

$$\tilde{t}_{x,y}(q) = \int_{-L/2}^{L/2} ds\, t_{x,y}(s) e^{-iqs},$$

we arrive at

$$V = V_0\zeta\frac{\omega^2}{\pi(\partial n/\partial\mu)}\int_0^\infty \frac{dp}{v^2} f_0' \,\mathrm{Im}\left[\tilde{t}_x\left(\frac{\omega}{v}\right)\tilde{t}_y\left(\frac{\omega}{v}\right)\right]. \quad (9)$$

For degenerate Fermi gas for which $\partial n/\partial\mu=2/\pi v_F$ it follows from Eq. (9) that

$$V = -V_0\zeta\frac{\omega^2}{2v_F^2}\mathrm{Im}\left[\tilde{t}_x\left(\frac{\omega}{v_F}\right)\tilde{t}_y\left(\frac{\omega}{v_F}\right)\right]. \quad (10)$$

We shall exemplify the general result (9) by means of the case of hard angle [Fig. 2(a)], where voltage takes the form

$$V = V_0\zeta\frac{1}{\pi(\partial n/\partial\mu)}\sin(2\alpha)\int_0^\infty dp\, f_0'[2\sin(\eta/2) - \sin\eta], \quad (11)$$

where $\eta=\omega L/v$. In a particular case of degenerate Fermi system Eq. (11) leads to

$$V = -V_0\zeta\frac{1}{2}\sin(2\alpha)[2\sin(\eta_F/2) - \sin\eta_F], \quad (\eta_F = \eta|_{v=v_F}). \quad (12)$$

Result (12) stays limited when $L\to\infty$, this implies that the voltage is formed by a finite curved part of the wire. Nevertheless, the remaining parts $s\sim L$ of the wire also participate in the voltage. The evidence of that is the presence of time-of-flight oscillations at frequencies $\omega_N=4N\pi v_F/L$ (N is integer) in Eq. (12) (which evidently survive for any considered shape of wires containing straight parts). The value of voltage, averaged with respect to oscillations, vanishes in the limit of large L. Just this value survives if the damping will be taken into account. The limited nature of result for hard angle leads to the finite contribution of the curved part.

IV. KINETIC APPROACH

Here we solve the kinetic equation (1) with collision term for infinitely long quantum wire, in assumption that equilibrium distribution is established inside the wire. The linear in $\mathcal{E}$ correction $f_1(t)=\mathrm{Re}(f_1^\omega e^{-i\omega t})$ to the equilibrium distribution function obeys to the equation

$$(-i\omega + iqv)\tilde{f}_1^\omega - e\tilde{\mathcal{E}}^\omega(q) v f_0' = -\nu[\tilde{f}_1^\omega - \langle\tilde{f}_1^\omega\rangle]. \quad (13)$$

The solution of Eq. (13) reads

$$\tilde{f}_1^\omega = -ie\tilde{\mathcal{E}}^\omega(q) v f_0' \frac{\omega + qv}{q^2v^2 - \omega^2 - i\omega\nu}. \quad (14)$$

For calculation of photocurrent we need stationary quadratic in $\mathbf{E}$ part of distribution function. It satisfies the equation

$$(\nu + iqv)\tilde{\tilde{f}}_2(p;q) - \frac{e}{4}\sum_{q'}\left[\tilde{\mathcal{E}}^{-\omega}(q-q')\frac{\partial\tilde{f}_1^\omega(p;q')}{\partial p} + (\omega\to-\omega)\right] = \nu\langle\tilde{\tilde{f}}_2(p;q)\rangle. \quad (15)$$

In an infinitely long wire the current tends to zero. This follows from infinite resistance of the system and finiteness of the size where the external field drags electrons. Hence the static drag current is compensated by the static Ohmic leakage current caused by appearing static electric field $\mathcal{E}_0$,

$$\tilde{j}(q) + \tilde{\sigma}(q)\tilde{\mathcal{E}}_0(q) = 0,$$

where $\tilde{\sigma}(q)$ is the linear conductivity. The photovoltage is

$$V \equiv \lim_{q\to 0} \tilde{\mathcal{E}}_0(q) = -\frac{\lim_{q\to 0} \tilde{j}(q)}{\tilde{\sigma}(0)}. \qquad (16)$$

Thus, the determination of photovoltage requires finding the limit of space Fourier component of photocurrent $\tilde{j}(q)$ at $q \to 0$. Using Eqs. (14) and (15) we obtain from Eq. (16)

$$V = -\frac{e^3}{(2\pi)^2\tilde{\sigma}(0)}\int_{-\infty}^{\infty} dp f_0' v^2 \frac{\partial(\tau v)}{\partial p} \times \int_{-\infty}^{\infty} dq q |\tilde{\mathcal{E}}^{\omega}(q)|^2 \frac{\omega\nu}{[(qv)^2-\omega^2]^2+(\omega\nu)^2}, \qquad (17)$$

where

$$\tilde{\sigma}(0) = -\frac{e^2}{\pi}\int_{-\infty}^{\infty} dp\, f_0' v^2 \tau.$$

With taking into account integration over q the value $|\tilde{\mathcal{E}}^{\omega}(q)|^2$ in Eq. (17) can be replaced with $2\,\mathrm{Im}(E_x^{\omega}E_y^{-\omega})\mathrm{Im}[\tilde{t}_x(q)\tilde{t}_y^*(q)]$.

At $\nu \to 0$ Eq. (17) is simplified to

$$V = -\frac{e^3}{4\pi\tilde{\sigma}(0)}\int_{-\infty}^{\infty} dp f_0' v^2 \frac{\partial(\tau v)}{\partial p} \times \int_{-\infty}^{\infty} dq q |\tilde{\mathcal{E}}^{\omega}(q)|^2 \delta[(qv)^2-\omega^2]. \qquad (18)$$

The physical meaning of Eq. (18) is very simple: in the presence of external alternating field the curvature induces package of waves $|\tilde{\mathcal{E}}^{\omega}(q)|^2$ any of which accelerates electrons moving with the velocity of wave $v=\omega/q$.

In the specific case of degenerate Fermi statistics Eq. (18) reads

$$V = -V_0\zeta\beta\frac{\omega^2}{2v_F^2}\mathrm{Im}\left[\tilde{t}_x\left(\frac{\omega}{v_F}\right)\tilde{t}_y\left(\frac{\omega}{v_F}\right)\right], \qquad (19)$$

where $\beta = 1+2\,\partial(\ln\tau)/\partial(\ln\varepsilon_F)$. If the relaxation time does not depend on the electron energy Eq. (19) reduces to the collisionless limit (10). Although both considerations give voltage independent from the scattering strength, a more accurate approach (19) depends on the character of scattering (via energy dependence of τ), despite the smallness of the scattering rate. In the case of weak scattering the mean-free path exceeds the length of curved domain and, thus, one should think that the scattering does not affect the voltage. The discrepancy can be explained by the fact that the voltage is formed on the same distance as the conductivity, namely, mean-free path $v\tau$ or, in other words, on the distance where the scattering occurs.

Let us consider another limit when the mean-free path is less than the length of the curved part. Neglecting q in the denominator of the ratio in Eq. (18) one can get the expression

$$V = -\frac{e^3}{2\pi\tilde{\sigma}(0)}\int_{-\infty}^{\infty} dp f_0' v^2 \frac{\tau}{\omega(\omega^2\tau^2+1)}\frac{\partial(\tau v)}{\partial p}\int_{-\infty}^{\infty} ds k(s), \qquad (20)$$

where $k(s)=[t_y'(s)t_x(s)-t_y(s)t_x'(s)]$ is the curvature. The integral $\int ds k(s)$ equals to the total angle of rotation of vector $\mathbf{t}$. This result physically follows from the locality of static field production in case of the small mean-free path: that means that the static field $\mathcal{E}(s)$ in the local case can be determined by nothing else but local curvature k. The universality of Eq. (20) gives immediately the same result for the cases of hard angle and open book and zero result for Ω curve (here we emphasize that all hard angles on curves are supposed to be smoother than l).

V. SPECIFIC SHAPES

When the mean-free path is comparable with the length of the curved part the induced voltage can be obtained from Eq. (17) by the substitution of specific expressions for $\tilde{t}_{x,y}(q)$, while Eqs. (18) and (19) refer to the limit of a large mean-free path.

For the hard angle $\tilde{\mathbf{t}}(q)$ has its singular behavior at $q=0$: $\tilde{t}_x(q)=2\pi\delta(q)\cos\alpha$ and $\tilde{t}_y(q)=(2i/q)\sin\alpha$. This singularity originates from behavior $\mathbf{t}(s)$ at $s\to\pm\infty$. Substituting $\tilde{\mathbf{t}}(q)$ to Eq. (17) we find

$$V = -V_0\zeta\beta\sin(2\alpha)\frac{\omega\tau}{1+(\omega\tau)^2}. \qquad (21)$$

In case of the open book [Fig. 2(b)] at $\nu\to 0$ and zero temperature Eq. (18) yields

$$V = -V_0\zeta\beta F(\xi), \qquad (22)$$

where

$$F(\xi) = \frac{2\xi - 2\xi\cos(2\alpha)\cos(2\alpha\xi) - (1+\xi^2)\sin(2\alpha)\sin(2\alpha\xi)}{2(1-\xi^2)^2}, \qquad (23)$$

$\pi-2\alpha$ is the angle at the vertex [see Fig. 2(b)], $\xi=R\omega/v_F$, and R is the radius of the circular part. If ξ is small or large (corresponding to small and large radius R), $F(\xi)\approx[1-\cos(2\alpha)-\alpha\sin(2\alpha)]\xi$ at $\xi\ll 1$ and $F(\xi)\approx -\sin(2\alpha)\sin(2\alpha\xi)/2\xi^2$ at $\xi\gg 1$.

Equation (18) and, hence, Eq. (23) are valid in the extreme case of $\nu\to 0$. If $R\to 0$, the open book converts to the hard angle and Eq. (23) should convert to Eq. (21) at $\omega\tau \gg 1$. This is not so. In fact, here we have a competition of two small parameters ξ and $1/\omega\tau$. Taking into account singular contributions to the Fourier transform $\tilde{\mathbf{t}}(q)$ (the same as in case of the hard angle) gives an additive contribution to voltage

$$- V_0 \zeta \beta \sin(2\alpha) \frac{1}{\omega\tau} \tag{24}$$

exactly coinciding with the hard angle result (21) at $\omega\tau \gg 1$.

For Ω-like wire the function $F(\xi)$ is replaced with

$$F(\xi) = \frac{2\xi^2}{(1-\xi^2)^2}\left[\xi \cos^2\left(\frac{\pi\xi}{2}\right) - \frac{1-\xi^2}{2}\sin(\pi\xi)\right]. \tag{25}$$

Note that in this case contribution (24) does not appear because straight parts of the curve are parallel. In limiting cases of large and small ξ, this function behaves as $F(\xi) \approx -(\pi - 2)\xi^3$ at $\xi \ll 1$ and $F(\xi) \approx \sin \pi\xi + (1+\cos \pi\xi)/\xi$ at $\xi \gg 1$.

VI. DISCUSSION

The considered ballistic and kinetic regimes differ in the participation of scattering in the voltage appearance. In the ballistic regime electrons go through the wire conserving the equilibrium states of the source and drain seas. In the kinetic regime the scattering and equilibrium establishment occurs inside the wire. This has no effect on the order of the voltage magnitude mainly determined by the parameter V_0 in both cases. Qualitatively, eV_0 is the mean kinetic energy which an electron obtains from the alternating field. The estimations give $V_0 \simeq 10^{-6}$ V for $E=1$ V/cm, $\omega = 10^{11}$ s^{-1}, and electron mass of GaAs $m = 0.07 \times 10^{-27}$ g.

The ballistic voltage experiences the time-of-flight oscillations. They smear out by scattering. In a particular case of hard angle this results in the disappearance of voltage in the kinetic regime. Under very strong scattering the static field becomes local as determined by the local curvature; hence, the global voltage is determined by a global geometry of the wire, namely, the total angle of rotation of the tangent. In the case of moderate scattering the voltage contains two contributions; one of which depends on the angle between straight entrance and exit and the other depends on the local geometry of the wire. The first of these contributions vanishes if the angle is π (for example, in Ω curve).

Note that in the kinetic regime the voltage goes to the finite limit if the length tends to infinity. This limit, as a rule, has the same order of magnitude as in the ballistic regime but, unlike the latter, is determined by the energy dependence of the relaxation. This difference is specified by the difference of relaxation: inside the wire in the kinetic case and outside the wire in the ballistic case.

The effect considered here occupies its place in a rank of other photoelectric effects including photon drag,[4–6] photogalvanic effect,[1–3] quantum pumps,[7–10] etc. It is desirable to compare them with each other. Unlike the photogalvanic effect which occurs in macroscopically uniform media, this curvature-induced effect has a local character. When the system size goes to infinity the produced voltage tends to the fixed limit instead of growth.

Let there are multiple kinks on the wire distributed with constant density. Then any kink will produce a fixed voltage, so that they additively produce the mean electric field (and the corresponding current in a short circuit regime). This variant reminds of the photocurrent in a spiral quantum wire.[15] Both effects have the same physical origin: the photon drag due to the curvature-induced effective momentum of the wave. The found voltage is independent on the relaxation rate in the limit of weak relaxation (except for the case of hard angle, where the absence of relaxation leads to the absence of voltage). At the same time the current in the spiral grows in the limit of weak relaxation under resonant conditions (when the velocity of wave coincides with the Fermi velocity). This difference between the effects results from the fact that the kink produces a wave package instead of a single wave in case of a spiral curve and averaging on momentums washes out the resonance and the dependence on the collision rate.

The system under consideration is akin to quantum pumps by locality of the alternating perturbation; but, unlike them, the found (with relaxation taken into account) current disappears in the infinite system. The asymmetry of considered systems is artificial. From this viewpoint these systems are similar to the artificial antidot lattices.[13,18,19] Being planar, the curved wires can be simply realized experimentally, in particular, in the same antidot lattices under depleting conditions. Hence, there is hope for a quick experimental verification.

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[1] V. I. Belinicher and B. I. Sturman, Usp. Fiz. Nauk **130**, 415 (1980) [Sov. Phys. Usp. **23**, 199 (1980)].

[2] E. L. Ivchenko and G. E. Pikus, in *Semiconductor Physics*, edited by V. M. Tuchkevich and V. Ya. Frenkel (Consultants Bureau, New York, 1986), p. 427.

[3] E. M. Basken, M. D. Blokh, M. V. Entin, and L. I. Magarill, Phys. Status Solidi B **83**, K97 (1977).

[4] A. M. Danishevskii, A. A. Kastal'skii, S. M. Ryvkin, and I. D. Varoshetskii, Zh. Eksp. Teor. Fiz. **58**, 544 (1970) [Sov. Phys. JETP **31**, 292 (1970)].

[5] A. A. Grinberg, Zh. Eksp. Teor. Fiz. **58**, 989 (1970) [Sov. Phys. JETP **31**, 531 (1970)].

[6] A. F. Gibson and S. Montasser, J. Phys. C **8**, 3147 (1975).

[7] D. J. Thouless, Phys. Rev. B **27**, 6083 (1983).

[8] P. W. Brouwer, Phys. Rev. B **58**, R10135 (1998).

[9] M. Moskalets and M. Büttiker, Phys. Rev. B **66**, 205320 (2002).

[10] L. S. Braginskii, M. M. Makhmudian, and M. V. Entin, JETP **100**, 920 (2005); M. M. Mahmoodian, L. S. Braginsky, and M. V. Entin, Phys. Rev. B **74**, 125317 (2006).

[11] R. D. Astumian and P. Hänggi, Phys. Today **55** (11), 33 (2002).

[12] P. Reimann, Phys. Rep. **361**, 57 (2002).

[13] A. D. Chepelianskii and D. L. Shepelyansky, Phys. Rev. B **71**,

052508 (2005).
[14] M. L. Leadbeater, C. L. Foden, J. H. Burroughes, M. Pepper, T. M. Burke, L. L. Wang, M. P. Grimphaw, and D. A. Ritchie, Phys. Rev. B **52**, R8629 (1995).
[15] L. I. Magarill and M. V. Entin, Pis'ma Zh. Eksp. Teor. Fiz. **78**, 249 (2003) [JETP Lett. **78**, 213 (2003)].
[16] M. V. Entin and M. M. Mahmoodian, EPL **84**, 47008 (2008).
[17] It should be emphasized that, from the point of view of quadratic response, a more general case of partly coherent radiation can be treated as a superposition of absolutely nonpolarized noncoherent light (which does not cause the current) and coherent monochromatic radiation, see, e.g., L. D. Landau and E. M. Lifshitz, *The Classical Theory of Fields*, Course of Theoretical Physics Vol. 2, 4th ed. (revised English edition) (Butterworth, Amsterdam, 1998), p. 133.
[18] A. D. Chepelianskii, M. V. Entin, L. I. Magarill, and D. L. Shepelyansky, Eur. Phys. J. B **56**, 323 (2007).
[19] S. Sassine, Yu. Krupko, J.-C. Portal, Z. D. Kvon, R. Murali, K. P. Martin, G. Hill, and A. D. Wieck, Phys. Rev. B **78**, 045431 (2008).