Dissipationless Electron Transport in Photon-Dressed Nanostructures

O.V. Kibis

Department of Applied and Theoretical Physics, Novosibirsk State Technical University,
Karl Marx Avenue 20, 630092 Novosibirsk, Russia

It is shown that the electron coupling to photons in field-dressed nanostructures can result in the ground electron-photon state with a nonzero electric current. Since the current is associated with the ground state, it flows without the Joule heating of the nanostructure and is nondissipative. Such a dissipationless electron transport can be realized in strongly coupled electron-photon systems with the broken time-reversal symmetry — particularly, in quantum rings and chiral nanostructures dressed by circularly polarized photons.

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The scientific trend, that has emerged in recent years, is to combine the methods of quantum optics with the advances in design and preparation of various nanostructures. This interdisciplinary and innovative field of physics, arisen from achievements of the modern nanotechnology, offers the possibility of both the fundamental studies of light-matter interaction in unusual artificial quantum systems and the development of optoelectronic nanodevices with unique characteristics (see, e.g., Ref. [1]). In the present Letter, the methodology of quantum optics is applied to analyze the strong interaction between electrons in curvilinear nanostructures and a quantized electromagnetic field. It follows from the analysis that the electron coupling to photons can result in the ground electron-photon state with a nonzero electric current. Since the current is associated to the ground state, it flows without the Joule heating of the nanostructure and is nondissipative. Thus, we declare a novel mechanism of the photon-induced superconductivity. The main goal of the Letter is to present the first theoretical analysis of this mechanism.

Let us consider the effect of an electromagnetic field on electron transport in curvilinear one-dimensional conductors (quantum wires) of a constant cross section and of a constant curvature — specifically, in ring-shaped and helix-shaped quantum wires (quantum rings and quantum helices, respectively). In what follows, the spin effects and the field-induced mixing of transverse electron modes in the quantum wire will be beyond consideration. Then, the processes of electron-field interaction, associated to the electron motion along the wire, can be described by the Hamiltonian \( \hat{H}_e = \frac{\hat{p}_s^2}{2m_e} + U(s) \), where \( s \) is the electron coordinate along the wire, \( \hat{p}_s = -i\hbar \partial/\partial s \) is the operator of electron momentum along the wire, \( U(s) = -e \int E_t(s)ds \) is the potential energy of an electron in the wire subjected to the electromagnetic field, \( e \) is the electron charge, \( m_e \) is the electron mass, and \( E_t(s) = \mathbf{E}t(s) \) is the projection of the electric field \( \mathbf{E} \) on a tangent to the quantum wire with the unit vector \( \mathbf{t}(s) \).

As to the magnetic component of the electromagnetic field, it does not influence the electron motion along the quantum wire since the Lorentz force is perpendicular to the wire. Considering the problem within the conventional quantum-field approach [2], the classical field \( \mathbf{E} \) should be replaced with the field operator \( \hat{\mathbf{E}} \). Then the full Hamiltonian of the electron-photon system, including both the field energy \( \hbar \omega \hat{a}^\dagger \hat{a} \) and the electron Hamiltonian \( \hat{H}_e \) is

\[
\hat{H} = \hbar \omega \hat{a}^\dagger \hat{a} + \frac{\hat{p}_s^2}{2m_e} - e \int \hat{\mathbf{E}}(s)ds ,
\]

where \( \hat{a} \) and \( \hat{a}^\dagger \) are the operators of photon annihilation and creation, respectively, written in the Schrödinger representation (the representation of occupation numbers) [2], and \( \omega \) is the frequency of the electromagnetic field. For definiteness, in this Letter we will discuss nanostructures exposed to a plane monochromatic circularly polarized electromagnetic wave. Assuming the electromagnetic wave to be clockwise polarized and neglecting any spatial inhomogeneity of the wave field, the field operator in Eq. [1] can be written as \( \hat{\mathbf{E}} = i \sqrt{2\pi \hbar \omega /V} (\mathbf{e}_+ \hat{a} - \mathbf{e}_- \hat{a}^\dagger) \), where \( V \) is the quantization volume, \( \mathbf{e}_\pm = (e_x \pm ie_y)/\sqrt{2} \) are the polarization vectors, \( e_x, y, z \) are the unit vectors directed along the \( x, y, z \) axes of the Cartesian coordinate system, and the \( z \) axis corresponds to the direction of the wave propagation [2].

First of all, let us consider a ring-shaped quantum wire (a quantum ring with the radius \( R \)) subjected to a circularly polarized electromagnetic wave propagating along the ring axis. Since the tangent vector to the ring is \( \mathbf{t}(s) = -\sin(s/R) \mathbf{e}_x + \cos(s/R) \mathbf{e}_y \), the Hamiltonian [1] takes the form \( \hat{H} = \hat{H}_R^{(0)} + \hat{U}_R \), where

\[
\hat{U}_R = -i e R \sqrt{\frac{\pi \hbar \omega}{V}} (e^{i\varphi} \hat{a} - e^{-i\varphi} \hat{a}^\dagger)
\]

is the Hamiltonian of the electron-photon interaction in the ring, \( \hat{H}_R^{(0)} = \hbar \omega \hat{a}^\dagger \hat{a} + \hbar^2 \hat{p}_s^2 / (2m_e R^2) \) is the Hamiltonian of a noninteracting electron-photon system, \( \hat{l}_z = -i \partial/\partial \varphi \) is the operator of the electron angular momentum along the ring axis, and \( \varphi = s/R \) is the angular coordinate in the ring. To describe the electron-photon system, let us use the notation \( |m, N\rangle \) which indicates that the
electromagnetic field is in a quantum state with the photon occupation number \( N = 1, 2, 3, \ldots \), and the electron is in a quantum state with the wave function \( \psi_m(\varphi) = \sqrt{1/2\pi} \exp(im\varphi) \), where \( m = 0, \pm 1, \pm 2, \ldots \) is the electron angular momentum along the ring axis. The electron-photon states \( |m, N\rangle \) are true eigenstates of the unperturbed Hamiltonian \( \hat{H}^{(0)} \) and their energy spectrum is \( \varepsilon_m^{(0)} = N\hbar\omega + \hbar^2 m^2/2m_e R^2 \). In order to find the energy spectrum of the full electron-photon Hamiltonian \( \hat{H} \), let us use the conventional perturbation theory, considering the term \( (2) \) as a perturbation with the matrix elements

\[
\langle m', N'|\hat{U}_R|m, N\rangle = -ieR\sqrt{\frac{\pi \hbar\omega}{\tau}} \left[ \sqrt{N}\delta_{m,m'-1}\delta_{N,N'+1} - \sqrt{N+1}\delta_{m,m'+1}\delta_{N,N'-1} \right].
\]

Performing trivial calculations within the second order of the perturbation theory, we arrive at the sought energy spectrum of the Hamiltonian \( \hat{H} \),

\[
\varepsilon_{m,N} = \varepsilon_m^{(0)} + \frac{\langle m+1,N-1|\hat{U}_R|m,N\rangle^2}{\varepsilon_m^{(0)} - \varepsilon_{m+1,N-1}^{(0)}} + \frac{\langle m-1,N+1|\hat{U}_R|m,N\rangle^2}{\varepsilon_m^{(0)} - \varepsilon_{m-1,N+1}^{(0)}}.
\]

Let us assume hereafter that the dressing electromagnetic field cannot be absorbed by electrons. In order to neglect the collisional absorption of the field, the condition

\[
\omega\tau \gg 1
\]

is assumed to be satisfied, where \( \tau \) is the effective electron life time restricted by a scattering of the electron from all defects — both geometrical and structural — of the nanostructure. To avoid the optical absorption, the field frequency \( \omega \) is assumed to be far from the resonant frequencies of the ring. Then the allowed energies of the complete electron-photon system are described by Eq. \( (3) \) for \( N = N_0 \), where \( N_0 \) is the constant photon occupation number of the field. If the electromagnetic wave is strong \( (N_0 \gg 1) \), the energy spectrum \( \varepsilon_m \) can be written as \( \varepsilon_{m,N_0} = N_0\hbar\omega + \varepsilon(m) \), where \( N_0\hbar\omega \) is the constant field energy, the electron-photon term is

\[
\varepsilon(m) = m^2 \varepsilon_R + eE_0R \left[ \frac{eE_0 R/2 \varepsilon_R}{(2m - \hbar\omega/\varepsilon_R)^2 - 1} \right],
\]

\( \varepsilon_R = \hbar^2/2m_e R^2 \) is the characteristic electron energy in the ring, \( E_0 = \sqrt{\pi N_0}\hbar\omega/V \) is the classical amplitude of electric field of the electromagnetic wave, and the factor in the square brackets is assumed to be much smaller than unity. As a consequence, we can consider electrons in the photon-dressed ring as quasiparticles with the energy spectrum \( \varepsilon_m \) which is shown schematically in Fig. 1(a) for the case of \( \hbar\omega/\varepsilon_R < 1 \). Let us restrict the analysis of the electric currents in the ring by the principal (zeroth) order of the perturbation theory. Then the current of a field-dressed electron with an angular momentum \( m \) is approximately equal to the current of a free electron with the same angular momentum, \( j_m = m e\hbar/2\pi R^2 m_e \) \[4\]. As a result, the current in the ring is

\[
\overline{j} = \sum_m j_m.
\]

Correspondingly, the full current in the ring is \( \overline{j} = 2j_0 + 2j_{-1} = -e\hbar/\pi R^2 m_e \). Assuming the radius of quantum ring \( R \) to be of the nanometer scale, we obtain \( j \sim 10^{-6} A \). Since this current corresponds to the ground state, it is nondissipative. To clarify the physical nature of the current, it should be reminded that the equality \( \varepsilon_{N} = \varepsilon_{-N} \) takes place in the absence of the field \( (E_0 = 0) \) due to the time-reversal symmetry (the Kramers theorem). Since the time reversal turns clockwise polarized photons into counterclockwise polarized ones and vice versa, the electron coupling to the circularly polarized field breaks this symmetry. Therefore, field-dressed electron states \( \psi_{m,N} \) with angular momenta \( m \) and \( -m \) are split, which results in the discussed effect. Since this splitting decreases with increasing \( m \), the effect can be observable in rings whose Fermi energy is sufficiently low (e.g., semiconductor rings).

Though the discussed nondissipative current differs substantially from the known persistent currents arising in quantum rings due to the Aharonov-Bohm effect \[4\], it still cannot be identified with the superconductivity declared in the preamble of the Letter. Indeed, this current flows in a microscopical ring and is devoid of such a character of superconductivity as a nondissipative translational motion of electrons. That is why persistent currents in quantum rings — regardless of their physical nature — are phenomena of purely academic interest. However, we will demonstrate hereafter that the same electron-photon coupling leads also to nondissipative electric currents in macroscopically long conductors. This argument is crucial to consider the discussed effect as a conceptually novel mechanism of superconductivity. Among the variety of different nanostructures, macroscopically long conductors with chiral symmetry — helicoidal quantum wires, chiral carbon nanotubes, DNA-based nanostructures, etc. — seem prospective for the realization of such a superconductivity. To substantiate this statement, we will discuss the simplest chiral nanostructure — a quantum helix (helic-shaped quantum wire). Such helices can be fabricated by methods of the modern nanotechnology \[6\] and promise a broad range of interesting phenomena \[7\].

Let us consider a right-hand quantum helix with a radius \( R \), a pitch \( P \), and a macroscopically large length
$L$. Then the tangent vector to the helix is $\mathbf{t}(s) = -gR\sin(gs)\mathbf{e}_x + gR\cos(gs)\mathbf{e}_y + (P/l)\mathbf{e}_z$, where $s$ is the coordinate along the helix, $l = \sqrt{(2\pi R)^2 + P^2}$ is the length of the helix turn, and $g = 2\pi/l$. Assuming a circularly polarized electromagnetic wave to be propagating along the axis of the helix, we can write the Hamiltonian (11) as $\mathcal{H} = \mathcal{H}^{(0)}_H + \hat{U}_H$, where

$$\hat{U}_H = -ieR\sqrt{\frac{\pi \hbar \omega}{V}} (e^{ig\mathbf{a} \cdot \mathbf{\hat{a}}} - e^{-ig\mathbf{\hat{a}} \cdot \mathbf{\hat{a}}}) \tag{6}$$

is the Hamiltonian of the electron-photon interaction in the helix, and $\mathcal{H}^{(0)}_H = \hbar \omega \mathbf{\hat{a}} \cdot \mathbf{\hat{a}} + p_e^2/2m_e$ is the Hamiltonian of a noninteracting electron-photon system. In order to neglect a disturbance of the field by the helix, the field frequency, $\omega$, is assumed to be within the transparency region of the helix [10]. To describe the electron-photon system, let us use the notation $| k, N \rangle$ which indicates that the electromagnetic field is in a quantum state with the photon occupation number $N = 1, 2, 3, \ldots$ and the electron is in a quantum state with the wave function $\psi_k(s) = \sqrt{1/L} \exp(iks)$, where $k$ is the electron wave vector along the helix. The electron-photon states $| k, N \rangle$ are true eigenstates of the unperturbed Hamiltonian $\mathcal{H}^{(0)}_H$ and their energy spectrum is $\varepsilon^{(0)}_{k,N} = N\hbar \omega + \hbar^2 k^2 / 2m_e$. Therefore, they form the complete basis of the electron-field interaction.

As a consequence, we arrive at the system of linear algebraic equations

$$\left(\varepsilon_{k,N} - \varepsilon_{k,N}^{(0)}\right) b_{k,N} + \langle k, N | \hat{U}_H | k + g, N - 1 \rangle b_{k+g,N-1} + \langle k, N | \hat{U}_H | k - g, N + 1 \rangle b_{k-g,N+1} = 0, \tag{8}$$

where $\varepsilon_{k,N}$ is the sought energy spectrum of the Hamiltonian (11). The equations (8) under the condition (11) describe the allowed energies of the electron-photon system $\varepsilon_{k,N}$ for the constant photon occupation number $N = N_0$. It follows from Eq. (7) that the electron-photon interaction (6) mixes a set of electron-photon states with wave vectors $k$ differing by $g$. If the characteristic energy of the electron-field interaction $eE_0R$ is much smaller than the characteristic energy interval between these states $\hbar^2 g^2 / 2m_e$, then it is sufficient to take into account the mixing of only two states closest in energy, neglecting the contribution of all other states. In this approximation the system of equations (8) is reduced to just two equations. Solving these two equations accurately for $N_0 \gg 1$, we obtain the energy spectrum of the electron-photon system, $\varepsilon_{k,N_0} = N_0 \hbar \omega + \varepsilon(k)$, where the energy of photon-dressed electron is

$$\varepsilon(k) = \begin{cases} \varepsilon_+(k) = \sqrt{\varepsilon^2(k) + (\Delta \varepsilon/2)^2} & \text{if } k - k_\omega < g/2, \\ \varepsilon_-(k) = \sqrt{\varepsilon^2(k) + (\Delta \varepsilon/2)^2} & \text{if } k - k_\omega > g/2, \end{cases} \tag{9}$$

$\varepsilon_\pm(k) = (\hbar^2 / 4m_e) (k^2 \pm k^2 / 2 + g^2) (k - k_\omega)$, $k_\omega = m_e \omega / g \hbar$, and $\Delta \varepsilon = \varepsilon | E_0 R |$. As a consequence, we can consider electrons in the field-dressed helix as quasiparticles with the energy spectrum (9) which is shown schematically in Fig. 1 (b). Correspondingly, the electric current of the quasiparticle with the wave vector $k$ is $j_k = ev_k / L$, where $v_k = (1 / \hbar) d\varepsilon(k) / dk$ is the quasiparticle velocity along the helix. Therefore the full current in the helix is $j = \sum_k j_k$, where the summation should be performed over filled states (9). As an example, the current of a field-dressed electron system with the Fermi energy $\mu$ (see Fig. 1 (b)) is $j = \int v_k dk / \pi = e \Delta \varepsilon / \pi \hbar$, where the integration is performed over the filled states lying under the Fermi level. Since this current is associated to the ground state of the macroscopically long photon-dressed quantum wire, it is nondissipative and corresponds to the decelerated photon-induced superconductivity.

![FIG. 1: (color online) The energy spectrum of electrons — in the quantum ring (a) and in the quantum helix (b) — dressed by the quantized circularly polarized field (heavy lines) and without the field (thin lines). The nondissipative currents $j$ correspond to the filled states (up) in the ring (a) and to the filled states lying under the Fermi energy $\mu$ in the helix (b).](image)
tering of electrons from such a periodic potential opens energy gaps at $k = \pm n g/2$ ($n = 1, 2, 3, \ldots$) in the electron energy spectra of various helicoidal nanostructures. In contrast to the stationary case, the rotating transverse electric field of the circularly polarized electromagnetic wave produces the periodic time-dependent potential, $U(s, t) = eE_0 R \cos(g s - \omega t)$, which runs along the helix with the velocity $v = \omega/g$. Correspondingly, the Bragg scattering of electrons from this running periodic potential takes place at $k' = \pm n g/2$, where $k' = k - k_\omega$ is the electron wave vector in the reference system moving along the helix with the same velocity $v$ (i.e., in the rest frame of the running potential). As a result, the Bragg gaps in the energy spectrum lie at the wave vectors, $k = k_\omega \pm g/2$, which correspond to the first diffraction maximum of the Bragg scattering ($n = 1$) and are positioned asymmetrically in the $k$ space of the laboratory reference system. As expected, in the limiting case of a stationary field ($\omega = 0$) this asymmetry vanishes ($k_\omega = 0$) and the spectrum exactly coincides with the electron energy spectrum of the quantum helix exposed to a stationary transverse electric field, which is given by Eq. (7) in Ref. 8. Certainly, the Bragg gaps also take place for $n = 2, 3, 4, \ldots$. However, these gaps are lying higher in energy, depend on higher powers of the electric field, $E_0$, and cannot be derived from the approximate expression. We will discuss them elsewhere.

It follows from the aforesaid that the Bragg scattering transfers electrons between the two states differing by the wave vector $g$ and the energy $\hbar \omega$ [see Fig. 1(b)]. It is also known that the conservation laws allow optical transitions of free electrons in the helix between the same two states. Therefore, the Bragg scattering can be reformulated in terms of the quantum optics: It is identical physically to the wave-induced Rabi oscillations of electrons between the mentioned two states with the Rabi frequency $\Omega = \Delta \varepsilon/\hbar$. Since collisional processes can erode any energy gap, the condition for the existence of the discussed gaps depends on the electron lifetime, $\tau$, and is

$$\Omega \tau \gg 1.$$  \hspace{1cm} (10)

As expected, the inequality coincides with the existence condition of the Rabi oscillations. Since the Rabi oscillations are not accompanied by any absorption of the field energy, the inequality can be interpreted as the forbidding of resonant field absorption at the wave frequency $\omega$. Therefore the inequalities and describe the purely dressing (nonabsorbable) field. Assuming the helix radius, $R$, to be of the nanometer scale, the energy gap is $\Delta \varepsilon \sim 10^{-2}$ eV for the field $E_0 \sim 10^5$ V/cm. In this case the nondissipative current is $j \sim 10^{-6}$ A.

It is crucial that the spectra and, which result in nondissipative currents, are derived for a nonabsorbable field satisfying the conditions and. Indeed, the absorption of the field energy by electrons leads to the usual photovoltaic (photon drag) effects in both quantum rings and quantum helices. Since such a field absorption is accompanied by the energy transfer from the field to electrons, photovoltaic currents result in the Joule heating and obey the Ohm law. Therefore, a dressing electromagnetic field producing nondissipative currents cannot be absorbable. We have demonstrated that the electron coupling to such a dressing field forms the physical mechanism of dissipationless electron transport, which differs substantially from the traditional mechanisms of the superconductivity. Generalizing the aforesaid, we can declare that any quantized periodic potential (a potential wave) running along a conductor can produce a nondissipative electric current, if the potential wave is both strong and of sufficiently high frequency. Formally, the quantized potential wave can be considered as a quantum version of the electron pump. The search for efficient ways to realize such a quantized pump in various conducting systems is the subject of future studies.

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* Electronic address: Oleg.Kibis@nstu.ru

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