Persistent current induced by quantum light

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It is demonstrated that the strong coupling of an electron gas to photons in systems with broken time-reversal symmetry results in bound electron-photon states which cannot be backscattered elastically. As a consequence, the electron gas can flow without dissipation. This quantum macroscopic phenomenon leads to the unconventional superconductivity which is analyzed theoretically for a two-dimensional electron system in a semiconductor quantum well exposed to an in-plane magnetic field.

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I. INTRODUCTION

Advances in laser physics achieved in recent decades have made possible the use of lasers as tools to manipulate the electronic properties of various quantum systems. Since the strong interaction between electrons and an intense laser field cannot be described as a weak perturbation, it is necessary to consider the system "electron + field" as a whole. Such a bound electron-photon object, which was called "electron dressed by field" (dressed electron), became commonly used model in modern physics.[1,2] The field-induced modification of the energy spectrum of dressed electrons—also known as a dynamic (ac) Stark effect—was discovered experimentally in both atoms[3] and solids[4] many years ago and has been studied in various electronic systems. Particularly, it is well known that the interaction between a solid and a monochromatic electormagnetic field can open energy gaps $\Delta\varepsilon$ within electron energy bands of the solid (see, e.g., Refs. 5–8). Such a gap opening is pictured schematically in Fig. 1(a) for the case of interaction between a bulk semiconductor with the band gap ε_g and an electromagnetic field with the frequency $\omega_0 > \varepsilon_g/\hbar$ (Ref. 6). It should be stressed that the gaps $\Delta\varepsilon$ are opened in the resonant points of **k** space (i.e., at electron wave vectors **k** satisfying the condition of "the photon energy $\hbar\omega_0$ is equal to the energy interval between electron bands"). If the electron energy spectrum of the solid is symmetric, $\varepsilon(\mathbf{k}) = \varepsilon(-\mathbf{k})$, the resonant points—and, correspondingly, the gaps $\Delta\varepsilon$—are positioned symmetrically in the **k** space with respect to band edges [see Fig. 1(a)]. Though the light-induced gap opening has been known for a long time,[5,6] its theory is developed exclusively for solids with such a symmetric electron energy spectrum. Electronic systems with an asymmetric energy spectrum $\varepsilon(\mathbf{k}) \neq \varepsilon(-\mathbf{k})$ escaped attention before and will be considered below.

It follows from the fundamentals of quantum mechanics that the asymmetric energy spectrum of electrons can exist in systems with broken both time-reversal symmetry and inversion symmetry. Particularly, it takes place in nanostructures without an inversion center in the presence of a magnetic field, including asymmetric quantum wells,[9–13] chiral carbon nanotubes,[14–16] hybrid semiconductor/ferromagnet nanostructures,[17] and so on. For definiteness, let us consider such a simple nanostructure as a quantum well (QW). Generally, a QW confines charge carriers, which were originally free to move in three dimensions (x,y,z), to two dimensions (x,y), forcing them to occupy a planar region inside the QW.[18] As a consequence, the physical properties of a QW depend on the electron potential energy $U(z)$ describing this confinement. In what follows, we will consider a QW with a confining potential $U(z)$ devoid of an inversion center (asymmetric QW). Technologically, such asymmetric QWs are fabricated on the basis of semiconductor heterojunctions (see, e.g., Refs. 18 and 19). If an asymmetric QW is exposed to an in-plane magnetic field H_y directed along the y axis, the electron energy spectrum of the QW consists of a set of subbands which are shifted along the k_x axis with respect to each other by the wave vector $\Delta k_x \propto H_y$ (Ref. 18). This shifting, which is schematically pictured in Fig. 1(b), leads to the asymmetric energy spectrum of electrons $\varepsilon(k_x) \neq \varepsilon(-k_x)$. Let an electron system with such an asymmetric energy spectrum be subjected to an electromagnetic field with the frequency ω_0. Then the resonant points of the intersubband electron-photon interaction are positioned asymmetrically in the **k** space with respect to subband edges. Correspondingly, it is reasonable to expect that the photon-induced energy gaps $\Delta\varepsilon$ will be positioned asymmetrically within the subbands. The remarkable feature of such an asymmetrically gapped energy spectrum is the nondissipative flowing of electron gas. For instance, let an electron gas fill dressed states under the Fermi level μ [see Fig. 1(b)]. It is easy to show that the electric current along the x axis, produced by the electron gas, is $j_x \propto \Delta\varepsilon$. Since this nonzero current is associated to the ground state of the electron system, it flows without dissipation. Thus, the photon-dressed electron system with broken time-reversal symmetry can demonstrate the superconductor-like behavior. The given paper is devoted to the theoretical justification of this phenomenon.

The paper is organized as follows. In the Sec. II we introduce the electron-photon Hamiltonian and find its exact eigenstates and eigenvalues. Section III is devoted to an analysis of electron transport in photon-dressed QWs. Section IV contains a discussion and conclusions.

II. ELECTRON-PHOTON HAMILTONIAN

The electron energy spectrum of a QW exposed to an in-plane magnetic field H_y can be described by the expression[18]

$$\varepsilon_n(\mathbf{k}) = \varepsilon_{n0} - \frac{\hbar\bar{z}_n e H_y}{cm^*}k_x + \frac{\hbar^2\left(k_x^2 + k_y^2\right)}{2m^*}, \qquad (1)$$

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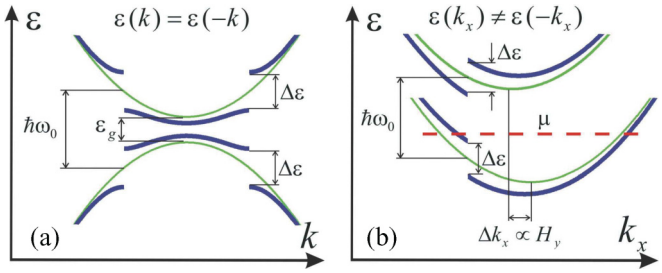

FIG. 1. (Color online) Energy spectrum of free electrons (thin lines) and electrons dressed by an electromagnetic field with the frequency ω_0 (solid lines): (a) conductivity band and valence band in a bulk semiconductor with the band gap ε_g; (b) first two electron subbands in an asymmetric quantum well exposed to an in-plane magnetic field H_y. The Fermi level of the electron system μ is pictured by the dashed line.

which takes into account linear magnetic-field terms. Here $\mathbf{k}$ is the in-plane electron wave vector, $n = 1,2,3,\ldots$, is the number of electron subbands, m^* is the effective electron mass in the QW, e is the electron charge, ε_{n0} is the energy of subband edge, and $\bar{z}_n = \langle\psi_n(z)|z|\psi_n(z)\rangle$ is the averaged z coordinate of the electrons in the QW. Correspondingly, electron wave functions in the subbands (1) are $\psi_n(\mathbf{k}) = \exp(i\mathbf{kr})\psi_n(z)$, where $\mathbf{r}$ is the in-plane radius vector, and the wave function $\psi_n(z)$, arising from the confining potential $U(z)$, meets the Schrödinger equation

$$-\frac{\hbar^2}{2m^*}\frac{d^2\psi_n(z)}{dz^2} + U(z)\psi_n(z) = \varepsilon_{n0}\psi_n(z).$$

Let us restrict our consideration by electron processes in the two lowest subbands (1) with $n = 1,2$. Hereafter we will mark these subbands $\varepsilon_{1,2}(\mathbf{k})$ and their wave functions $\psi_{1,2}(\mathbf{k})$ by the symbols $\varepsilon^{\pm}(\mathbf{k})$ and $\psi^{\pm}(\mathbf{k})$, respectively, where the sign "−" corresponds to the lower subband in Fig. 1(b) ($n = 1$), and the sign "+" corresponds to the upper one ($n = 2$). Correspondingly, the shifting of the electron subbands in $\mathbf{k}$ space, pictured in Fig. 1(b), is $\Delta k_x = (\bar{z}_1 - \bar{z}_2)eH_y/c\hbar$.

Let the QW be exposed to a plane monochromatic electromagnetic wave with the frequency ω_0 (dressing field), which is linearly polarized along the z axis. For simplicity, we will neglect any spatial inhomogeneity of the dressing field, but a proper generalization can be easily made. To describe the wave-induced mixing of the electron states from the two subbands $\varepsilon^+(\mathbf{k})$ and $\varepsilon^-(\mathbf{k})$, the electron Hamiltonian should be written as a 2×2 matrix $\hat{\mathcal{H}}_e = [\varepsilon^+(\mathbf{k}) + \varepsilon^-(\mathbf{k})]\hat{I}/2 + [\varepsilon^+(\mathbf{k}) - \varepsilon^-(\mathbf{k})]\hat{\sigma}_z/2$, where $\hat{I}$ is the unity matrix, and $\hat{\sigma}_{x,y,z}$ are the Pauli matrices written in the basis of the two electron states $\psi^{\pm}(\mathbf{k})$. As to the Hamiltonian of the intersubband electron-wave interaction, it can be expressed within the dipole approximation by $\hat{\mathcal{H}}_{\rm int} = -dE\hat{\sigma}_x$, where $\mathbf{E} = (0, 0, E)$ is the electric field vector of the wave (the dressing field vector), and $d = e\langle\psi_1(z)|z|\psi_2(z)\rangle$ is the intersubband matrix element of the electric dipole moment which is assumed to be real and positive. Considering the problem within the conventional quantum-field approach,[1,2] the classical field E should be replaced with the field operator $\hat{E} = i\sqrt{2\pi\hbar\omega_0/V}(\hat{a} - \hat{a}^\dagger)$, where V is the quantization volume of the field, $\hat{a}$ and $\hat{a}^\dagger$ are the photon operators of annihilation and creation, respectively, written in the Schrödinger picture (the representation of occupation numbers[20]). After such a replacement, the interaction Hamiltonian takes the form $\hat{\mathcal{H}}_{\rm int} = -id\sqrt{2\pi\hbar\omega_0/V}(\hat{\sigma}_+\hat{a} + \hat{\sigma}_-\hat{a} - \hat{\sigma}_-\hat{a}^\dagger - \hat{\sigma}_+\hat{a}^\dagger)$, where $\hat{\sigma}_\pm = (\hat{\sigma}_x \pm i\hat{\sigma}_y)/2$. In what follows, we will assume that the photon energy $\hbar\omega_0$ is much more than the characteristic energy of the electron-field interaction dE_0, where E_0 is the field amplitude. Then the terms $\hat{\sigma}_+\hat{a}^\dagger$ and $\hat{\sigma}_-\hat{a}$ can be neglected, which corresponds to the rotating-wave approximation commonly used in quantum optics.[1,2] As a result, the interaction Hamiltonian is $\hat{\mathcal{H}}_{\rm int} = -id\sqrt{2\pi\hbar\omega_0/V}(\hat{\sigma}_+\hat{a} - \hat{\sigma}_-\hat{a}^\dagger)$. Since the operator of the field energy is $\hat{\mathcal{H}}_0 = \hbar\omega_0\hat{a}^\dagger\hat{a}$, the full Hamiltonian of the considered electron-photon system, $\hat{\mathcal{H}} = \hat{\mathcal{H}}_0 + \hat{\mathcal{H}}_e + \hat{\mathcal{H}}_{\rm int}$, reads as

$$\hat{\mathcal{H}} = \hbar\omega_0\hat{a}^\dagger\hat{a} + \frac{\varepsilon^+(\mathbf{k}) + \varepsilon^-(\mathbf{k})}{2}\hat{I} + \frac{\varepsilon^+(\mathbf{k}) - \varepsilon^-(\mathbf{k})}{2}\hat{\sigma}_z - id\sqrt{\frac{2\pi\hbar\omega_0}{V}}(\hat{\sigma}_+\hat{a} - \hat{\sigma}_-\hat{a}^\dagger). \quad (2)$$

The Hamiltonian (2) is formally similar to the Hamiltonian of the exactly solvable Jaynes-Cummings model.[21] Therefore, the Schrödinger problem with the electron-photon Hamiltonian (2) can also be solved exactly. Applying the methodology[22,31] to the problem, let us introduce the joined electron-photon space $|\pm,N\rangle = |\psi^{\pm}(\mathbf{k})\rangle \otimes |N\rangle$ to describe the electron being in the state with the wave function $\psi^{\pm}(\mathbf{k})$ and the dressing field being in the state with the photon occupation number $N = 1,2,3,\ldots$. The basic states of this space $|\pm,N\rangle$ are orthonormal and meet the conditions $\langle\pm,N|\pm,N'\rangle = \delta_{N,N'}$ and $\langle\pm,N|\mp,N'\rangle = 0$. Therefore, the exact eigenstates of the Hamiltonian (2), $\varphi_N^+(\mathbf{k})$ and $\varphi_N^-(\mathbf{k})$, can be written as

$$|\varphi_N^{\pm}(\mathbf{k})\rangle = \sqrt{\frac{\Omega_\pm(\mathbf{k}) + |\omega(\mathbf{k})|}{2\,\Omega_\pm(\mathbf{k})}}\,|\pm,N\rangle + i\eta(\mathbf{k})\sqrt{\frac{\Omega_\pm(\mathbf{k}) - |\omega(\mathbf{k})|}{2\,\Omega_\pm(\mathbf{k})}}\,|\mp,N\pm 1\rangle, \quad (3)$$

where $\Omega_\pm(\mathbf{k}) = \sqrt{8d^2(N + 1/2 \pm 1/2)(\pi\omega_0/\hbar V) + \omega^2(\mathbf{k})}$, $\omega(\mathbf{k}) = \omega_0 - [\varepsilon^+(\mathbf{k}) - \varepsilon^-(\mathbf{k})]/\hbar$, and

$$\eta(\mathbf{k}) = \begin{cases} -1, & \omega(\mathbf{k}) > 0, \\ 1, & \omega(\mathbf{k}) \leqslant 0. \end{cases}$$

The two energy branches corresponding to the electron-photon states (3), $\varepsilon_N^+(\mathbf{k})$ and $\varepsilon_N^-(\mathbf{k})$, are given by

$$\varepsilon_N^{\pm}(\mathbf{k}) = N\hbar\omega_0 + \frac{\varepsilon^+(\mathbf{k}) + \varepsilon^-(\mathbf{k})}{2} \pm \frac{\hbar\omega_0}{2} \pm \eta(\mathbf{k})\frac{\hbar\Omega_\pm(\mathbf{k})}{2}. \quad (4)$$

The expressions (3) and (4) can be easily verified by direct substitution into the Schrödinger equation $\hat{\mathcal{H}}\varphi_N^{\pm}(\mathbf{k}) = \varepsilon_N^{\pm}(\mathbf{k})\varphi_N^{\pm}(\mathbf{k})$ with the Hamiltonian (2), keeping in mind the relations[20,23]

$$\hat{\sigma}_\pm\,|\mp,N\rangle = |\pm,N\rangle\,, \quad \hat{\sigma}_\pm\,|\pm,N\rangle = 0\,,$$
$$\hat{a}\,|\pm,N\rangle = \sqrt{N}\,|\pm,N-1\rangle, \quad \hat{\sigma}_z\,|\pm,N\rangle = \pm|\pm,N\rangle,$$
$$\hat{a}^\dagger\,|\pm,N\rangle = \sqrt{N+1}\,|\pm,N+1\rangle.$$

As expected, the state $|\varphi_N^{\pm}(\mathbf{k})\rangle$ turns into the state $|\pm,N\rangle$ when the electron-photon interaction vanishes (i.e., when $d = 0$). In the most interesting case of a laser-generated intense dressing

field (when the values of N and V tend to infinity while the ratio N/V is constant[1]), the full energy of the electron-photon system (4) can be written as a sum $N\hbar\omega_0 + \varepsilon(\mathbf{k})$, where $N\hbar\omega_0$ is the energy of the dressing field, and

$$\varepsilon(\mathbf{k}) = \frac{\varepsilon^+(\mathbf{k}) + \varepsilon^-(\mathbf{k})}{2} \pm \frac{\hbar\omega_0}{2} \pm \eta(\mathbf{k})\frac{\Omega(\mathbf{k})}{2} \tag{5}$$

is the desired energy spectrum of the dressed electrons. Here $\Omega(\mathbf{k}) = \sqrt{(\hbar\Omega_R)^2 + [\hbar\omega_0 - \varepsilon^+(\mathbf{k}) + \varepsilon^-(\mathbf{k})]^2}$, $\Omega_R = dE_0/\hbar$ is the Rabi frequency of intersubband electron transitions, and $E_0 = \sqrt{8\pi N\hbar\omega_0/V}$ is the classical amplitude of the dressing field. The two branches of the spectrum (5) describe two subbands of dressed electrons, which are pictured in Fig. 1(b) with solid lines: The branch with the sign "+" corresponds to the upper pictured subband and the branch with the sign "−" corresponds to the lowest one. As expected, the energy gaps

$$\Delta\varepsilon = \hbar\Omega_R, \tag{6}$$

pictured in Fig. 1(b), are opened at the electron wave vector $\mathbf{k} = \mathbf{k}_0$ satisfying the resonance condition $\hbar\omega_0 = \varepsilon^+(\mathbf{k_0}) - \varepsilon^-(\mathbf{k_0})$. If the dressing field is absent ($E_0 = 0$), the expression (5) turns into the the spectrum of the free electron, $\varepsilon_{1,2}(\mathbf{k})$, given by Eq. (1) and pictured in Fig. 1(b) with thin lines. It should be noted that the energy spectrum of dressed electrons, written in the form (5), is of universal character and applicable to describe dressed electron states arising from any two photon-mixed energy bands of free electrons $\varepsilon^-(\mathbf{k})$ and $\varepsilon^+(\mathbf{k})$ in any solid. Particularly, in the case of a direct gap semiconductor, the energy spectrum of the free electrons has the form $\varepsilon^\pm(\mathbf{k}) = \pm\varepsilon_g/2 \pm \hbar^2\mathbf{k}^2/2m^*$, where ε_g is the semiconductor band gap. Substituting this expression into Eq. (5), we arrive at the energy spectrum of dressed electrons in a bulk semiconductor, which is pictured in Fig. 1(a) and was shown for the first time in Ref. 6.

III. TRANSPORT PROPERTIES OF DRESSED ELECTRONS

The electric current along the x axis, which is produced by a dressed electron with the wave vector $\mathbf{k}$, is $j_x(\mathbf{k}) = ev_x(\mathbf{k})/L_x$, where $L_{x,y}$ are the in-plane dimensions of the QW along the x,y axes, and $\mathbf{v}(\mathbf{k})$ is the average velocity of the dressed electron in the state (3). This average velocity is given by the conventional quantum-mechanical expression $\mathbf{v}(\mathbf{k}) = \langle\varphi_N^\pm(\mathbf{k})|\hat{\mathbf{v}}|\varphi_N^\pm(\mathbf{k})\rangle$, where $\hat{\mathbf{v}} = (i/\hbar)[\hat{\mathcal{H}}\hat{\mathbf{r}} - \hat{\mathbf{r}}\hat{\mathcal{H}}]$ is the operator of the electron velocity, and $\hat{\mathbf{r}}$ is the operator of the electron coordinate. To calculate the matrix element $\langle\varphi_N^\pm(\mathbf{k})|\hat{\mathbf{v}}|\varphi_N^\pm(\mathbf{k})\rangle$, we will use the $\mathbf{k}$ representation of the coordinate operator $\hat{\mathbf{r}} = i\partial/\partial\mathbf{k}$ (Ref. 23). Then the expression for the velocity of the dressed electron takes the form $\mathbf{v}(\mathbf{k}) = (1/\hbar)\partial\varepsilon_N^\pm(\mathbf{k})/\partial\mathbf{k}$. As expected, this expression coincides formally with the well-known classical Hamilton equation $\mathbf{v}(\mathbf{p}) = \partial\varepsilon(\mathbf{p})/\partial\mathbf{p}$, describing the velocity of a particle with an energy $\varepsilon(\mathbf{p})$ and the generalized particle momentum $\mathbf{p} = \hbar\mathbf{k}$. In the case of the many-electron system, the full current along the x axis is $j_x = \sum_\mathbf{k} j_x(\mathbf{k})$, where the summation should be performed over filled states (3). Taking into account the aforesaid, the current produced by dressed electrons lying under the Fermi level μ [see Fig. 1(b)] is $j_{x0} = (eL_y/2\pi^2\hbar)\int(\partial\varepsilon(\mathbf{k})/\partial k_x)d^2\mathbf{k}$, where the integration should be performed over the spectrum (5) with the energy $\varepsilon(\mathbf{k}) \leqslant \mu$. As a result of the integration, we arrive at the expression

$$j_{x0} = \frac{eL_y}{\pi^2\hbar^2}\sqrt{2m^*\mu}\,\Delta\varepsilon. \tag{7}$$

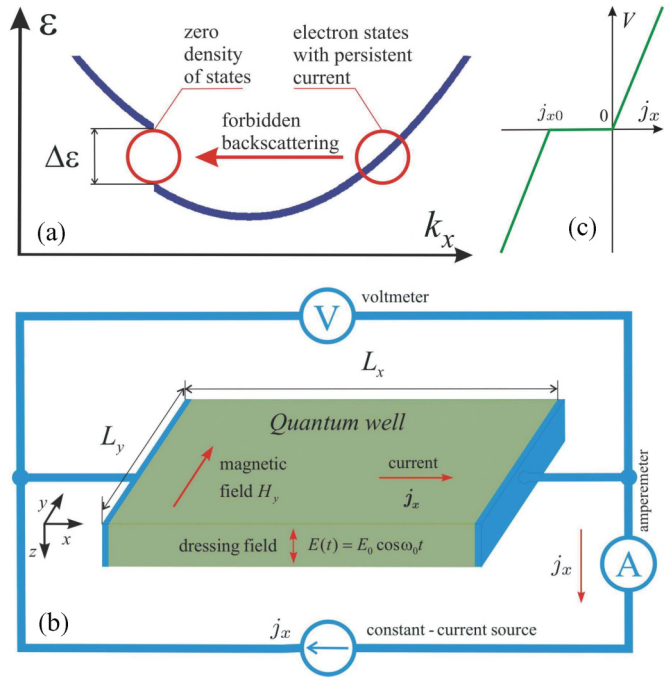

FIG. 2. (Color online) Scheme of the light-induced superconductivity: (a) asymmetrically gapped energy spectrum of electrons as a physical reason of the persistent current; (b) sketch of a proposed experimental setup for electron transport measurements in a photon-dressed quantum well; (c) voltage-current characteristic of the photon-dressed quantum well.

Since the current (7) is associated to the ground state of the photon-dressed electron system, it flows without dissipation and should be considered as a persistent current. It follows from Eq. (7) that the current arises from the photon-induced energy gap (6): If the dressing field is absent ($E_0 = 0$), both the gap (6) and the current (7) vanish. The physical reason of the interrelationship between the persistent current (7) and the gap (6) is clarified in Fig. 2(a). Since the gap (6) forbids an elastic backscattering of electrons from states lying in the energy range $\Delta\varepsilon$, a current associated to these electron states flows without dissipation. Therefore, we mark these states in Fig. 2(a) as "electron states with persistent current." As expected, the full current of the marked states is exactly equal to the persistent current (7). It should be stressed that the gap (6) arises from stationary solutions of the time-independent Schrödinger problem with the stationary Hamiltonian (2) describing the closed system "electron + quantized electromagnetic field." Therefore, it is the true (stationary) gap in the density of electron-photon states (3). As a consequence, the gap (6) will manifest itself directly in all phenomena sensitive to the density of states of charge carriers, including backscattering processes marked by the arrow in Fig. 2(a).

It follows from the charge conservation law that an electric current in any conductor must satisfy the continuity condition.

Applying this general rule to the considered system, we arrive at the evident result: The ground state of a QW with the nonzero current (7) satisfies the continuity condition only if the QW is a part of a closed electrical circuit, where the current $j_x = j_{x0}$ flows [see Fig. 2(b)]. If the circuit is broken, the continuity condition requires the zero current $j_x = 0$. Particularly, a state of the photon-dressed electron system with the zero current always takes place in an isolated QW: Though the state is not ground, it corresponds to the minimal energy under the additional condition of $j_x = 0$. As a consequence, the expected voltage-current characteristic of the QW, $V(j_x)$, takes the form pictured schematically in Fig. 2(c). The remarkable feature of the characteristic is the zero electrical resistance in the broad range of currents $j_{x0} < j_x < 0$, which arises from the electron states with persistent current. Though such a zero-resistance behavior of photon-dressed QW is superconductor-like, the physical reason of the declared effect differs conceptually from both the conventional superconductivity in solids and superfluidity in quantum liquids. Indeed, these known mechanisms of the nondissipative flow are based on interaction processes in strongly correlated many-particle systems, whereas the discussed phenomenon arises from the one-electron Hamiltonian (2) and takes place for a photon-dressed gas of noninteracting electrons. Thus, the declared effect results in an unconventional superconductivity which is discussed below.

IV. DISCUSSION AND CONCLUSION

The electron-photon Hamiltonian (2) does not take into account a scattering of electrons, which leads to the finite lifetime of electron states (1), τ. Therefore, the results obtained above from the Hamiltonian (2) are correct if the photon-induced energy gap $\Delta\varepsilon$ is much large than the scattering-induced washing of the electron energy spectrum $\hbar/\tau$. Taking into account Eq. (6), this condition of applicability of the Hamiltonian (2) can be written in the form

$$\Omega_R \tau \gg 1, \tag{8}$$

which coincides with the condition of photon-induced Rabi oscillations of electrons between the subbands $\varepsilon^-(\mathbf{k})$ and $\varepsilon^+(\mathbf{k})$ (Ref. 2). Thus, the gap (6)—and, correspondingly, the light-induced persistent current (7)—takes place when the electron subsystem is in the regime of intersubband Rabi oscillations. The coexistence of the persistent current and the Rabi oscillations has a deep physical meaning. To clarify it, let us write the eigenstates (3) in the resonant point $\mathbf{k} = \mathbf{k}_0$, where the gap (6) is opened. Since the wave vector $\mathbf{k}_0$ satisfies the condition $\omega(\mathbf{k}_0) = 0$, the eigenstates (3) at $\mathbf{k} = \mathbf{k}_0$ are

$$|\varphi_N^{\pm}(\mathbf{k}_0)\rangle = \frac{1}{\sqrt{2}}\left[|\pm, N\rangle + i\,|\mp, N \pm 1\rangle\right], \tag{9}$$

and their energies (5) are $\varepsilon^{\pm}(\mathbf{k}_0) \pm \Delta\varepsilon/2$, respectively. It is seen that the bound electron-photon states (9) describe the one-photon mixing of the electron subbands $\varepsilon^-(\mathbf{k})$ and $\varepsilon^+(\mathbf{k})$, which corresponds physically to the Rabi oscillations of the electron subsystem between these two subbands. As to the gap (6), it can be treated as a total binding energy of the two coherent electron-photon states (9) with the two different signs "±". Thus, the existence of the photon-induced coherent states (9) with the binding energy (6) is microscopical reason of the discussed effect.

It should be noted that an electron, performing periodical Rabi oscillations between the subbands $\varepsilon^-(\mathbf{k})$ and $\varepsilon^+(\mathbf{k})$, does not absorb the energy of electromagnetic field inducing these Rabi oscillations. In other words, the energy of the dressing field averaged over the period of Rabi oscillations is constant. As a consequence, the condition of Rabi oscillations (8) is physically equal to the forbidding of all processes accompanied with the absorption of the field energy by electrons. Therefore, the inequality (8) can be considered as a condition of a purely dressing (nonabsorbable) electromagnetic field. As a result, the existence of a persistent current under the condition (8) does not contradict the energy conservation law since the nondissipative flow of photon-dressed electrons is not accompanied with the absorbing of energy of the dressing field.

It follows from the aforesaid that the light-induced persistent current (7) differs conceptually from light-induced currents which arise from a photovoltaic effect. Indeed, any photovoltaic effect must be accompanied with absorbing light energy by electrons (see, e.g., Ref. 24) and, therefore, it is absent under the condition (8). On the contrary, if the condition (8) is broken, the intersubband absorption of light—and, correspondingly, the photovoltaic effect—appears. In other words, there are two different regimes of electron-photon interaction in systems with broken time-reversal symmetry: (I) The regime of weak electron-photon coupling ($\Omega_R\tau \lesssim 1$), where an usual ohmic current appears from the photovoltaic effect and (II) the regime of strong electron-photon coupling ($\Omega_R\tau \gg 1$), where the persistent current (7) exists. The first regime has been studied in QWs both theoretically and experimentally,[25–29] whereas the second one escaped attention before. It should be stressed that the light-induced ohmic current and the light-induced persistent current, which arise from the two different physical mechanisms, flow in mutually opposite directions. Therefore, they can be easily differentiated in experiments.

Since the synthesis of semiconductor QWs with such characteristic parameters as an energy interval between electron subbands $[\varepsilon^+(\mathbf{k}_0) - \varepsilon^-(\mathbf{k}_0)] \sim 10^{-1}$ eV, electron lifetime $\tau \sim 10^{-10}$ s, intersubband dipole moment $d \sim 10^2$ D, and Fermi energy $\mu \sim 10^{-2}$ eV is routine procedure for modern nanotechnology, the appropriate source of a dressing field for the experimental observation of the declared effect seems an infrared laser—for instance, an ordinary CO_2-laser with the wavelength $\lambda = 10.6\ \mu$m and the output power $P \gtrsim 10$ W/cm^2. It should be noted that infrared lasers have been discussed as prospective sources of a dressing field for observing dressed electron states in semiconductors.[30] In our case, such a laser satisfies both the condition of the resonant electron-photon interaction $\hbar\omega_0 = \varepsilon^+(\mathbf{k}_0) - \varepsilon^-(\mathbf{k}_0)$ and the condition (8). The current (7) induced by this laser field can be estimated for the above-mentioned parameters as $j_{x0}/L_y \gtrsim 10^{-2}$ A/m. Thus, the persistent current (7) is large enough to be observed experimentally in actual semiconductor nanostructures.

Though we calculated the persistent current (7) for the electron distribution corresponding to zero temperature, the discussed effect can also exist at nonzero temperatures.

Indeed, the asymmetrically gapped energy spectrum pictured in Fig. 2(a) leads always to a nonzero current if an electron distribution function depends only on electron energy. The particular case of such a distribution function is the Fermi-Dirac function describing an electron distribution on dressed states (5) in the thermodynamic equilibrium at any temperature. Since there is no dissipation processes in the thermodynamic equilibrium, the discussed persistent current takes place at nonzero temperatures as well. In other words, the discussed superconductivity can be high temperature: It exists unless the increasing of electron scattering—which always takes place with increasing temperature—washes the gap (6).

Generally, the light-induced persistent current needs broken time-reversal symmetry. In the current paper we have considered an electron system which is devoid of time-reversal symmetry due to a magnetic field. However, the similar effect can also take place for a time-reversally symmetric electron system interacting with a dressing field without time-reversal symmetry. For instance, a circularly polarized field is devoid of time-reversal symmetry since the time reversal turns clockwise polarized photons into counterclockwise polarized ones and vice versa. Therefore, the electron coupling to a circularly polarized field can result in the persistent current in various time-reversally symmetric electron systems—particularly, in curvilinear quantum wires.[31] However, a fabrication of the quantum wires is not a trivial technological problem that impedes an experimental observation of the discussed effect in one-dimensional conductors. On the contrary, semiconductor QWs with high-mobility two-dimensional electron gas can be easily synthesized in a modern laboratory. As a result, the significant new area of experimental research in QWs—where the physics of superconductivity, the physics of nanostructures, and quantum optics meet—can be opened.

It should be noted that there is another kind of persistent current associated with the ground state of an electron system without time-reversal symmetry. This known persistent current arises from the Aharonov-Bohm effect[32] and takes place in quantum rings exposed to a magnetic field.[33,34] However, the Aharonov-Bohm persistent current cannot be identified with superconductivity since this current is closed. Indeed, it flows in a microscopical ring and is devoid of such a characteristic property of superconductivity as a dissipationless carrying of electrons over a macroscopically long distance. In contrast to the Aharonov-Bohm persistent current, the light-induced persistent current (7) flows in a QW with macroscopically large in-plane dimensions $L_{x,y}$ [see Fig. 2(b)]. Therefore, the persistent current (7) leads to the dissipationless carrying of an electric charge over the macroscopically long distance L_x. This allows to consider the discussed effect as a conceptually novel mechanism of superconductivity.

Finalizing the discussion, it should be noted that the appearance of unusual effects of strong electron-photon coupling is a common feature of quantum systems with broken fundamental symmetries, including both the broken time-reversal symmetry[31] and the broken inversion symmetry.[35] In the current paper we have solved the spinless problem, where the considered effect arises from a diamagnetic transformation of the electron energy spectrum (1) by a magnetic field H_y. However, the similar effect can also appear in photon-dressed spin systems without time-reversal symmetry. Since an analysis of spin-originated effects goes beyond the scope of the current paper, it will be done elsewhere.

In summary, we have declared the nondissipative flowing of photon-dressed electron gas (the light-induced superconductivity), which differs conceptually from both the superfluidity of quantum liquids and conventional superconductivity in solids. This macroscopic quantum phenomenon lies in the scientific area where condensed-matter physics and quantum optics meet. It is of universal character and can take place in various strongly coupled electron-photon systems with broken time-reversal symmetry. Particularly, the phenomenon can be observed in asymmetric quantum wells exposed to an in-plane magnetic field. Since the fabrication of such quantum wells is routine procedure for modern nanotechnology, the declared phenomenon can take place in actual nanostructures.

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