Metal-insulator transition in graphene induced by circularly polarized photons

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Exact stationary solutions of the electron-photon Dirac equation are obtained to describe the strong interaction between massless Dirac fermions in graphene and circularly polarized photons. It follows from them that this interaction forms bound electron-photon states which should be considered as a kind of charged quasiparticles. The energy spectrum of the quasiparticles is of dielectric type and characterized by an energy gap between the valence and conductivity bands. Therefore the electron-photon interaction results in metal-insulator transition in graphene. The stationary energy gap, induced by photons, and concomitant effects can be observed for graphene exposed to a laser-generated circularly polarized electromagnetic wave.

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I. INTRODUCTION

Since the discovery of graphene[1]—a monolayer of carbon atoms—its unique physical characteristics have aroused enormous interest in the scientific community. Particularly, the influence of electromagnetic field on electronic properties of graphene is in focus of attention.[2] However all theoretical models, elaborated to describe this influence, deal with strong electromagnetic field in the framework of classical electrodynamics. A theory capable of describing the strong interaction between electrons in graphene and quantized electromagnetic field was unknown up to the present. The paper is aimed to fill partially this gap in the theory. It is surprising that the problem of interaction between electrons at the Fermi level of graphene and circularly polarized photons can be solved exactly in analytical form without the need of any approximation and numerical calculations. This is one of few exactly solvable problems in the theory of photon-matter interaction that is not without interest to both condensed-matter physics and quantum optics. Moreover, since electrons in graphene behave as massless relativistic fermions,[1,3,4] graphene can serve as a low-energy "proving ground" for the relativistic physics.[5] Therefore the obtained solutions can also be useful to describe the interaction between relativistic particles and photons. Thus the solved problem lies at the intersection of different excited fields of modern physics that is what caused to write this paper.

II. DIRAC PROBLEM FOR ELECTRON-PHOTON SYSTEM IN GRAPHENE

Generally, electron states in graphene near the Fermi energy are described by eight-component wave functions written in a basis corresponding to two crystal sublattices of graphene, two electron valleys, and two orientations of electron spin.[2] In what follows intervalley scattering processes and spin effects will be beyond consideration, that reduces the number of necessary wave-function components to two. To describe the interaction between a plane monochromatic circularly polarized electromagnetic wave and electrons in a graphene, we shall be to use the Cartesian coordinate system (x,y,z), where the axis z is perpendicular to the graphene fixed at $z=0$. Then near the point at the Fermi level, where the valence and conductivity bands of graphene touch each other (the Dirac point), the Hamiltonian of electrons interacting with the electromagnetic field can be written for a single valley and for a certain direction of electron spin in the form[2]

$$\hat{\mathcal{H}}_e = v_F \hat{\boldsymbol{\sigma}} \left(\hat{\mathbf{p}} - \frac{e}{c}\mathbf{A} \right), \tag{1}$$

where v_F is Fermi velocity, $\hat{\mathbf{p}}$ is operator of electron momentum in the graphene plane, e is electron charge, and $\mathbf{A}$ is vector potential of the electromagnetic field. As to the vector operator $\hat{\boldsymbol{\sigma}}$, its components are Pauli matrices written in the basis of two orthogonal electron states arisen from two crystal sublattices of graphene. Since the operator $\hat{\sigma}_z$ is diagonal, these two states can be denoted by mutually opposite orientations of the pseudospin along the z axis, $s_z = \pm 1/2$. Considering the problem within the standard quantum-field approach,[6] the classical field, $\mathbf{A}$, should be replaced with the field operator, $\hat{\mathbf{A}}$. Assuming the electromagnetic wave to be clockwise polarized and propagating along the z axis, this operator can be written as

$$\hat{\mathbf{A}} = \sqrt{2\pi\hbar c^2/\omega_0 V}(\mathbf{e}_+ \hat{a} + \mathbf{e}_- \hat{a}^\dagger), \tag{2}$$

where ω_0 is frequency of the electromagnetic wave, V is volume of space with the field, $\mathbf{e}_\pm = (\mathbf{e}_x \pm i\mathbf{e}_y)/\sqrt{2}$ are polarization vectors, $\mathbf{e}_{x,y}$ are unit vectors directed along the x,y axes, $\hat{a}$ and $\hat{a}^\dagger$ are photon operators of annihilation and creation, respectively, written in the Schrödinger representation (the representation of occupation numbers).[6] Then the complete Hamiltonian of the electron-photon system, including both the field energy, $\hbar\omega_0 \hat{a}^\dagger \hat{a}$, and the electron Hamiltonian (1) rewritten in the quantum-field form is given by

$$\hat{\mathcal{H}} = \hbar\omega_0 \hat{a}^\dagger \hat{a} + v_F \hat{\boldsymbol{\sigma}} \hat{\mathbf{p}} - e\sqrt{\frac{4\pi\hbar v_F^2}{\omega_0 V}}(\hat{\sigma}_+ \hat{a} + \hat{\sigma}_- \hat{a}^\dagger), \tag{3}$$

where $\hat{\sigma}_\pm = (\hat{\sigma}_x \pm i\hat{\sigma}_y)/2$ are step-up and step-down operators for the z projection of the pseudospin.

The stationary solutions of the effective Dirac equation with the Hamiltonian (3) have the form $\Psi(\mathbf{k})=e^{i\mathbf{kr}}\psi(\mathbf{k})$, where $\mathbf{k}$ is wave vector in the graphene plane, $\mathbf{r}$ is radius vector, and $\psi(\mathbf{k})$ is eigenstate of the Hamiltonian

$$\hat{\mathcal{H}}_{\mathbf{k}}=\hbar\omega_0\hat{a}^{\dagger}\hat{a}+\hbar v_F\hat{\boldsymbol{\sigma}}\mathbf{k}-e\sqrt{\frac{4\pi\hbar v_F^2}{\omega_0 V}}(\hat{\sigma}_{+}\hat{a}+\hat{\sigma}_{-}\hat{a}^{\dagger}). \quad (4)$$

At the Dirac point ($\mathbf{k}=0$) the Hamiltonian (4) is formally similar to the Hamiltonian of exactly solvable Jaynes-Cummings model.[7] As a consequence, the effective Dirac equation with the Hamiltonian (4) for $\mathbf{k}=0$ can also be solved exactly. To describe the electron-photon system at the Dirac point, let us use the notation $|s_z,N\rangle$ which indicates that the electron is in one of two quantum states with the pseudospin projections $s_z=\pm 1/2$ and the electromagnetic field is in quantum state with the photon occupation number $N=1,2,3,\ldots$. Then exact eigenstates of the Hamiltonian (4) for $\mathbf{k}=0$, $\psi(0)=\varphi_{+1/2,N}$, and $\psi(0)=\varphi_{-1/2,N}$ can be written as

$$|\varphi_{\pm 1/2,N}\rangle=\sqrt{\frac{\Omega_{\pm}+\omega_0}{2\Omega_{\pm}}}|\pm 1/2,N\rangle \pm\frac{e}{|e|}\sqrt{\frac{\Omega_{\pm}-\omega_0}{2\Omega_{\pm}}}|\mp 1/2,N\pm 1\rangle, \quad (5)$$

where

$$\Omega_{\pm}=\sqrt{16(N+1/2\pm 1/2)(\pi v_F^2 e^2/\hbar\omega_0 V)+\omega_0^2},$$

energies of the electron-photon states [Eq. (5)], $\varepsilon_{+1/2,N}$ and $\varepsilon_{-1/2,N}$, are given by

$$\varepsilon_{\pm 1/2,N}=N\hbar\omega_0\pm\frac{\hbar\omega_0}{2}\mp\frac{\hbar\Omega_{\pm}}{2}, \quad (6)$$

and subscript indexes in Eqs. (5) and (6) indicate genesis of these states: the state $|\varphi_{\pm 1/2,N}\rangle$ turns into the state $|\pm 1/2,N\rangle$ when the electron-photon interaction vanishes (i.e., for $e=0$). Expressions (5) and (6) can be easily verified by direct substitution into the effective Dirac equation $\hat{\mathcal{H}}_{\mathbf{k}}\varphi_{\pm 1/2,N}=\varepsilon_{\pm 1/2,N}\varphi_{\pm 1/2,N}$ with the Hamiltonian (4) for $\mathbf{k}=0$, keeping in mind the trivial relations[6,8]

$$\hat{\sigma}_{\pm}|\mp 1/2,N\rangle=|\pm 1/2,N\rangle,\quad \hat{\sigma}_{\pm}|\pm 1/2,N\rangle=0,$$

$$\hat{a}^{\dagger}|\pm 1/2,N\rangle=\sqrt{N+1}|\pm 1/2,N+1\rangle,$$

$$\hat{a}|\pm 1/2,N\rangle=\sqrt{N}|\pm 1/2,N-1\rangle.$$

Let the photon occupation number of the electromagnetic field (2) is fixed at $N=N_0$ by a field source. Then the electron-photon states $\varphi_{-1/2,N_0}$ and $\varphi_{+1/2,N_0}$, originated from electron states degenerated at the Dirac point, are separated by the energy gap $\varepsilon_g=\varepsilon_{-1/2,N_0}-\varepsilon_{+1/2,N_0}$. In what follows we shall be to assume the electromagnetic wave to be classically strong, that corresponds to macroscopically large photon occupation numbers ($N_0\gg 1$). Then the gap can be written as

$$\varepsilon_g=\sqrt{W_0^2+(\hbar\omega_0)^2}-\hbar\omega_0, \quad (7)$$

where $W_0=2v_F eE_0/\omega_0$ is energy of electron rotational motion induced by the circularly polarized wave and $E_0=\sqrt{4\pi N_0\hbar\omega_0/V}$ is classical amplitude of electric field of the wave.

Let us proceed to solving the effective Dirac equation with the Hamiltonian (4) near the Dirac point for wave vectors $\mathbf{k}\neq 0$. Since the obtained electron-photon states [Eq. (5)] are eigenstates of the same Hamiltonian (4) for $\mathbf{k}=0$, they form complete basis of the considered electron-photon system. Therefore true eigenstates of the Hamiltonian (4) for $\mathbf{k}\neq 0$ can be sought as an expansion

$$\psi(\mathbf{k})=\sum_{m=1}^{\infty}[c_{+1/2,m}(\mathbf{k})\varphi_{+1/2,m}+c_{-1/2,m}(\mathbf{k})\varphi_{-1/2,m}]. \quad (8)$$

Substituting expansion (8) into the effective Dirac equation $\hat{\mathcal{H}}_{\mathbf{k}}\psi(\mathbf{k})=\varepsilon(\mathbf{k})\psi(\mathbf{k})$ with the Hamiltonian (4), we arrive at the system of recurrent algebraic equations for coefficients $c_{\pm 1/2,m}(\mathbf{k})$. This system can be easily solved in the special case when the effective parameter of electron-field interaction, $\alpha=W_0/\hbar\omega_0$, satisfies the condition $\alpha\ll 1$. Seeking the dispersion dependencies $\varepsilon(\mathbf{k})$ and $c_{\pm 1/2,m}(\mathbf{k})$ for any $\mathbf{k}$ in the principal order with respect to α, the system of recurrent equations can be reduced to the two equations

$$\left(\varepsilon-N_0\hbar\omega_0-\frac{\varepsilon_g}{2}\right)c_{-1/2,N_0}=\hbar v_F(k_x+ik_y)c_{+1/2,N_0}, \quad (9)$$

$$\left(\varepsilon-N_0\hbar\omega_0+\frac{\varepsilon_g}{2}\right)c_{+1/2,N_0}=\hbar v_F(k_x-ik_y)c_{-1/2,N_0}, \quad (10)$$

which can be solved exactly. Solving Eqs. (9) and (10) by conventional methods, we obtain two branches of the energy spectrum, $\varepsilon(\mathbf{k})$, which can be written as

$$\varepsilon_{N_0}^{\pm}(\mathbf{k})=N_0\hbar\omega_0\pm\sqrt{(\varepsilon_g/2)^2+(\hbar v_F k)^2}, \quad (11)$$

where omitted $\mathbf{k}$ terms are $\sim o(\alpha)$. As a result, eigenstates of the Hamiltonian (3) are given by

$$\Psi_{N_0}^{\pm}(\mathbf{k})=e^{i\mathbf{kr}}[c_{+1/2,N_0}^{\pm}(\mathbf{k})\varphi_{+1/2,N_0}+c_{-1/2,N_0}^{\pm}(\mathbf{k})\varphi_{-1/2,N_0}], \quad (12)$$

where superscript indices "$\pm$" denote solutions of Eqs. (9) and (10) corresponding to the two energy branches [Eq. (11)]. As expected, for $\mathbf{k}=0$ the obtained expressions (11) and (12) coincide with the exact solutions (5) and (6).

III. ELECTRON-PHOTON QUASIPARTICLES IN GRAPHENE AND THEIR PROPERTIES

The full energy of the electron-photon system (11) consists of the field energy, $N_0\hbar\omega_0$, and the term

$$\Delta\varepsilon(\mathbf{k})=\pm\sqrt{(\varepsilon_g/2)^2+(\hbar v_F k)^2}, \quad (13)$$

which is arisen from electron-photon interaction. As a consequence, an electron interacting with circularly polarized photons can be considered formally as a quasiparticle with the energy spectrum [Eq. (13)]. From the viewpoint of classical electrodynamics, the circularly polarized electromagnetic wave rotates the electron.[9] Therefore the quasiparticle

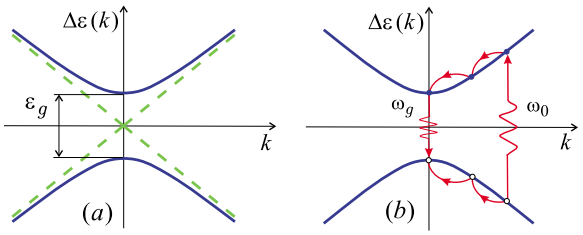

FIG. 1. (Color online) (a) Energy spectrum of quasiparticles in graphene in the presence of circularly polarized irradiation (solid line) and in the absence of one (dashed line); (b) scheme of quasiparticle transitions in graphene irradiated at the frequency ω_0.

is the mix of photon states and electron states corresponding to this rotational motion of the electron. Certainly, a center of trajectory of the rotating electron can move in the graphene plane that corresponds to a translational motion of the quasiparticle as a whole. The translational motion is described in Eq. (13) by the wave vector of quasiparticle, $\mathbf{k}$. Since the quasiparticle is *rotating* electron dressed by circularly polarized photons, this wave vector differs from the wave vector of free electron in graphene.

The energy spectrum of quasiparticles [Eq. (13)], schematically pictured in Fig. 1(a), defines optical and transport properties of graphene irradiated by circularly polarized light. The distinctive feature of the spectrum [Eq. (13)] is the energy gap [Eq. (7)] between quasiparticle states originated from valence and conductivity bands of graphene. Formally, the value of the photon-induced gap in the limit of large radiation intensities, given by Eq. (7), is the same as the value of dynamical gap induced at the Dirac point of graphene by a classical circularly polarized electromagnetic wave.[10,11] However, there is a conceptual difference between the *stationary* energy gap [Eq. (7)] arisen from quantum-field theory and *dynamical* gaps[10–13] arisen in graphene from classical fields. Historically, the term "dynamical gap" is come into the theory of graphene-field interaction[10–13] from the general theory[14,15] developed to describe different quantum systems exposed to time-dependent classical fields and denotes the gap in spectrum of Floquet quasienergies. Thus the dynamical gap is not true gap in the density of states of charge carriers. As it was specially stressed in Ref. 12, "since it is time-dependent problem, one cannot measure the gap directly from the density of states." In contrast to the dynamical gap, the gap [Eq. (7)] is obtained from stationary solutions of the time-independent Dirac problem for the electron-photon system and is true gap in the density of bound electron-photon states (charged quasiparticles). As a consequence, the stationary gap [Eq. (7)] will be to manifest itself directly in all phenomena sensitive to the density of states of charge carriers. In other words, the incorporation of such an additional physical factor as a quantum nature of electromagnetic field into the theory of graphene-field interaction, presented in the given paper, allows to consider the strong interacting electron-photon system in graphene as a gas of noninteracting quasiparticles with density of states defined by the energy spectrum [Eq. (13)]. Since the quasiparticle concept is conventional approach of the modern physics to describe quantum systems with a strong interaction, the developed theory is fruitful to predict and analyze different unexplored phenomena in graphene irradiated by light—particularly, physical effects discussed hereafter.

If another electromagnetic field with the frequency ω is applied to the considered electron-photon system, optical transitions of quasiparticles between its valence and conductivity bands can be possible. It follows from the energy conservation law that the transitions are allowed for $\omega \geq \omega_g$, where $\omega_g = \varepsilon_g/\hbar$. Therefore the gap [Eq. (7)] results in the threshold of absorption of an external field at the frequency $\omega = \omega_g$. Since the threshold frequency, ω_g, is controlled by the field intensity, E_0, this effect can be used as a basis for tunable detectors of electromagnetic radiation. The transition rate at the threshold frequency is defined by the interband dipole matrix element

$$|\langle \Psi^+_{N_0}(0)|\hat{d}_i|\Psi^-_{N_0}(0)\rangle| = |e|\hbar v_F \frac{\varepsilon_g + 2\hbar\omega_0}{2\varepsilon_g(\varepsilon_g + \hbar\omega_0)}, \qquad (14)$$

where $\hat{\mathbf{d}}$ is dipole moment operator and $i = x, y$. Besides optical transitions induced by the external field, there is the quasiparticle transition at the frequency ω_0,

$$\Psi^-_{N_0}(\mathbf{k}) \rightarrow \Psi^+_{N_0-1}(\mathbf{k}), \qquad (15)$$

which is accompanied by one-photon absorption of the field [Eq. (2)]. This transition, pictured in Fig. 1(b), needs a special discussion since the field [Eq. (2)] is inherent in the solved Dirac problem and cannot be considered as an external field. It follows from Eq. (11) that for $\omega_g \leq \omega_0$ there is the wave vector, $\mathbf{k}'$, satisfying the condition $\varepsilon^-_{N_0}(\mathbf{k}') = \varepsilon^+_{N_0-1}(\mathbf{k}')$. Therefore the transition [Eq. (15)] is allowed by the energy conservation law for $\omega_g \leq \omega_0$ and can take place at $\mathbf{k} = \mathbf{k}'$ without the assistance of external energy sources. However the field [Eq. (2)] has been accounted exactly by the complete Hamiltonian of the electron-photon system [Eq. (3)] and cannot lead directly to transitions between different stationary states of the same Hamiltonian. To realize the transition [Eq. (15)], the eigenstates of the Hamiltonian (3), $\Psi^-_{N_0}(\mathbf{k})$ and $\Psi^+_{N_0-1}(\mathbf{k})$, should be mixed by an external elastic scatterer (for example, by a defect of graphene lattice). Using Eq. (12), we can calculate the matrix element

$$\langle \Psi^+_{N_0-1}(\mathbf{k})|U(\mathbf{r})|\Psi^-_{N_0}(\mathbf{k})\rangle, \qquad (16)$$

which defines the rate of transition [Eq. (15)] for any scattering field $U(\mathbf{r})$. In the absence of scatterers, when $U(\mathbf{r}) = 0$ [or, generally, $U(\mathbf{r}) = \text{const}$], the matrix element [Eq. (16)] is zero and the transition is impossible. Thus the quasiparticle transition [Eq. (15)] is scatterer-assisted process.

It is seen in Fig. 1(b) that excited quasiparticles, accomplishing a cascade of nonradiative thermalization transitions after the transition [Eq. (15)], can create a population inversion between the conductivity band and the valence band at $\mathbf{k} = 0$. Since the interband dipole moment [Eq. (14)] is nonzero that results in light emission from graphene at the threshold frequency, ω_g, controlled by the field intensity, E_0.

Therefore the scheme, pictured in Fig. 1(b), forms a physical basis for creating tunable sources of electromagnetic radiation. The proposed mechanisms for tunable detection and generation of electromagnetic radiation are especially challenging for the terahertz (THz) frequency range. Indeed, search for effective THz sources and detectors is one of most excited problems of modern applied physics.[16,17] Moreover, the latest trend is the using of different nanostructures to fill the THz gap.[18–22] Therefore the application of graphene for THz devices fits well current tendencies in the graphene-based quantum electronics[23] and the nanophotonics as a whole.

Formally, the energy spectrum [Eq. (13)] is of dielectric type for any field intensity, E_0. However the interband transition [Eq. (15)] leads to generation of free-charge carriers in valence and conductivity bands for $\omega_g \leq \omega_0$. To turn graphene into true insulator, we need to forbid the transition [Eq. (15)] by the energy conservation law. This forbidding corresponds to $\omega_g > \omega_0$. Therefore the metal-insulator transition in graphene is a threshold effect and occurs at the critical field intensity, $E_0 = E_c$, satisfying the condition $\varepsilon_g = \hbar\omega_0$. Using Eq. (7), we obtain

$$E_c = \frac{\sqrt{3}\hbar\omega_0^2}{2v_F|e|}. \tag{17}$$

Being in the insulator state for $E_0 > E_c$, graphene has zero conductivity for the temperature $T=0$ and behaves as a semiconductor for $T \neq 0$. Applying the conventional theory of semiconductors[24] to graphene for $E_0 > E_c$ and $T \ll \varepsilon_g$, we can write the total density of free charge carriers for a single valley of graphene and for a certain direction of electron spin in the form

$$n = \left(\frac{m^* T}{\pi\hbar^2}\right)\exp\left(-\frac{\varepsilon_g}{2T}\right), \tag{18}$$

where m^* is effective mass of quasiparticles, which defines their energy spectrum,

$$\Delta\varepsilon(\mathbf{k}) = \pm\varepsilon_g/2 \pm \hbar^2\mathbf{k}^2/2m^*, \tag{19}$$

for small quasiparticle wave vectors $\mathbf{k}$. Considering $\mathbf{k}$ terms in the Hamiltonian (4) as a perturbation and using the standard perturbation theory,[8] we can find the energy spectrum of quasiparticles near their band edge as a $\mathbf{k}$-power series expansion without assuming weakness of the field (the parameter of electron-field interaction, $\alpha = W_0/\hbar\omega_0$, is allowed to be not small). Obtaining terms $\sim\mathbf{k}^2$ in the spectrum, we arrive at the exact expression for the effective mass

$$m^* = \frac{2\varepsilon_g(\varepsilon_g + 2\hbar\omega_0)(\varepsilon_g + \hbar\omega_0)^2}{v_F^2[(\varepsilon_g + 2\hbar\omega_0)^3 + \varepsilon_g^3]}. \tag{20}$$

It should be noted that the energy spectrums of quasiparticles [Eqs. (13) and (19)] are mutually complementary: Eq. (13) is applicable within a broad range of wave vectors but only for small field intensities ($\alpha \ll 1$) and Eq. (19) with the effective mass [Eq. (20)] is valid for any field intensities but only for small wave vectors ($\hbar v_F k/\varepsilon_g \ll 1$). As expected, Eqs. (13) and (19) are equal for small field intensities and small wave vectors.

IV. DISCUSSION AND CONCLUSIONS

Finalizing the paper, we have to estimate the effects discussed above. The distance between neighboring electron-photon states [Eq. (6)] is characterized by the two energies, $\hbar\omega_0$ and $\varepsilon_g = \hbar\omega_g$. Therefore the destructive influence of scattering processes on quasiparticles can be neglected if $\omega_0\tau \gg 1$ and $\omega_g\tau \gg 1$, where $\tau \sim 10^{-12}c$ is mean free time of charge carriers in graphene.[2] Thus the developed concept of quasiparticles is applicable if the field (2) is both high frequency and strong. Particularly, such a field can be created by using lasers. Let the wave is generated by a laser with the wavelength λ and focused into a narrow beam with the diameter $d \approx \lambda$. Then the wave amplitude is $E_0 \approx 4\sqrt{P/c\lambda^2}$, where P is output power of the laser. Within this scheme the condition of metal-insulator transition [Eq. (17)] can be satisfied, for instance, by using an ordinary low-power CO_2 laser with $\lambda = 10.6$ μm and $P \approx 76$ Watts. In this case the gap [Eq. (7)] is $\varepsilon_g \approx 117$ meV that allows to observe the metal-insulator transition at room temperatures. As to the effect of THz emission from irradiated graphene, it takes place for substantially smaller powers $P \sim 10^{-2}$ W.

As it follows from the estimations given above, the considered phenomena can be observable in graphene under a low-power irradiation. This is surprising since a *high-power* laser radiation is usually necessary to observe effects of strong electron-photon coupling.[25,26] Let us give simple physical reasons to clarify this unexpected result. From the quasiclassical viewpoint, the characteristic energy of electron coupling to circularly polarized electromagnetic wave is the kinetic energy of electron rotational motion induced by the wave. In the case of an usual condensed matter with parabolic electron energy spectrum, $\varepsilon(k) = \hbar^2k^2/2m^*$, this kinetic energy is $W_0^* = (eE_0/\omega_0)^2/2m^*$. As to graphene, due to its linear electron energy spectrum, $\varepsilon(k) = \pm\hbar v_F|k|$, the energy of the electron rotational motion induced by the wave is $W_0 = 2v_F eE_0/\omega_0$. Since we have $W_0/W_0^* \gg 1$ for weak fields E_0, the electron coupling to a low-power circularly polarized electromagnetic radiation is substantially stronger in graphene than in other condensed matters. The observability of the discussed effects for a low-power laser pumping is crucial from experimental viewpoint since high-power lasers fluidize a crystal lattice and are inapplicable to condensed-matter experiments. Thus graphene gives an unique opportunity to observe effects of strong electron-photon coupling, which cannot be observed in other condensed matters. As a result, the presented theory opens a significant new area of graphene-related research, where condensed-matter physics and quantum optics meet.

We have demonstrated that the Dirac electron spectrum of graphene leads to the strong electron interaction with circularly polarized photons for low light intensities. This results

in bound electron-photon states (charged quasiparticles) which should be considered as a substantially new kind of field-matter coupling. The energy spectrum of the quasiparticles is of dielectric type that leads to the photon-induced metal-insulator transition. From applied viewpoint, this effect can be used as a basis for creating tunable optoelectronic devices.

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[1] K. S. Novoselov, A. K. Geim, S. V. Morozov, D. Jiang, Y. Zhang, S. V. Dubonos, I. V. Grigorieva, and A. A. Firsov, Science **306**, 666 (2004).
[2] A. H. Castro Neto, F. Guinea, N. M. R. Peres, K. S. Novoselov, and A. K. Geim, Rev. Mod. Phys. **81**, 109 (2009), and references therein.
[3] P. K. Wallace, Phys. Rev. **71**, 622 (1947).
[4] I. A. Luk'yanchuk and Y. Kopelevich, Phys. Rev. Lett. **93**, 166402 (2004).
[5] G. W. Semenoff, Phys. Rev. Lett. **53**, 2449 (1984).
[6] V. B. Berestetskii, E. M. Lifshitz, and L. P. Pitaevskii, *Quantum Electrodynamics* (Pergamon Press, Oxford, 1982).
[7] E. T. Jaynes and F. W. Cummings, Proc. IEEE **51**, 89 (1963).
[8] L. D. Landau and E. M. Lifshitz, *Quantum Mechanics: Nonrelativistic Theory* (Pergamon Press, Oxford, 1991).
[9] L. D. Landau and E. M. Lifshitz, *The Classical Theory of Fields* (Pergamon Press, Oxford, 1987).
[10] T. Oka and H. Aoki, Phys. Rev. B **79**, 081406(R) (2009).
[11] T. Oka and H. Aoki, J. Phys.: Conf. Ser. **148**, 012061 (2009).
[12] F. J. López-Rodríguez and G. G. Naumis, Phys. Rev. B **78**, 201406(R) (2008); **79**, 049901(E) (2009).
[13] S. V. Syzranov, M. V. Fistul, and K. B. Efetov, Phys. Rev. B **78**, 045407 (2008).
[14] T. Dittrich, P. Hänggi, G. Ingold, B. Kramer, G. Schön, and W. Zwerger, *Quantum Transport and Dissipation* (Wiley, New York, 1998) and references therein.
[15] S. Kohler, J. Lehmann, and P. Hänggi, Phys. Rep. **406**, 379 (2005), and references therein.
[16] B. Ferguson and X. C. Zhang, Nature Mater. **1**, 26 (2002).
[17] M. Lee and M. C. Wanke, Science **316**, 64 (2007).
[18] D. Dragoman and M. Dragoman, Prog. Quantum Electron. **28**, 1 (2004).
[19] O. V. Kibis, M. Rosenau da Costa, and M. E. Portnoi, Nano Lett. **7**, 3414 (2007).
[20] M. E. Portnoi, O. V. Kibis, and M. Rosenau da Costa, Superlattices Microstruct. **43**, 399 (2008).
[21] K. G. Batrakov, S. A. Maksimenko, P. P. Kuzhir, and C. Thomsen, Phys. Rev. B **79**, 125408 (2009).
[22] O. V. Kibis, G. Ya. Slepyan, S. A. Maksimenko, and A. Hoffmann, Phys. Rev. Lett. **102**, 023601 (2009).
[23] D. Dragoman and M. Dragoman, Prog. Quantum Electron. **33**, 165 (2009).
[24] A. I. Anselm, *Introduction to Semiconductor Theory* (Prentice-Hall, Englewood Cliffs, NJ, 1981).
[25] M. O. Scully and M. S. Zubairy, *Quantum Optics* (University Press, Cambridge, 2001).
[26] C. Cohen-Tannoudji, J. Dupont-Roc, and G. Grynberg, *Atom-Photon Interactions: Basic Processes and Applications* (Wiley, Weinheim, 2004).