Electromagnetic-Field-Induced Suppression of Transport through *n-p* Junctions in Graphene

M. V. Fistul[1] and K. B. Efetov[1,2]
[1]*Theoretische Physik III, Ruhr-Universität Bochum, D-44801 Bochum, Germany*
[2]*L. D. Landau Institute for Theoretical Physics, 117940 Moscow, Russia*
(Received 20 March 2007; published 20 June 2007)

We study electronic transport through an *n-p* junction in graphene irradiated by an electromagnetic field (EF). In the absence of EF one may expect the perfect transmission of quasiparticles flowing perpendicular to the junction. We show that the resonant interaction of propagating quasiparticles with the EF induces a *dynamic gap* between electron and hole bands in the quasiparticle spectrum of graphene. In this case the strongly suppressed quasiparticle transmission is only possible due to interband tunneling. The effect may be used to control transport properties of diverse structures in graphene, e.g., *n-p-n* transistors and quantum dots, by variation of the intensity and frequency of the external radiation.

DOI: 10.1103/PhysRevLett.98.256803 PACS numbers: 73.63.−b, 05.60.Gg, 73.43.Jn, 81.05.Uw

Recent success in fabrication of graphene samples [1,2] has resulted in a stream of publications devoted to the study of this interesting material. The unique properties of graphene originate from peculiarities of the electron spectrum. The quasiparticle spectrum $\epsilon(p)$ consists of two valleys, and in each valley there are electron and hole bands crossing each other at some point. Near these points the electron spectrum is linear,

$$\epsilon_{\pm}(p) = \pm v|\mathbf{p}|, \tag{1}$$

where $\vec{\mathbf{p}} = \{p_x, p_y\}$ is the quasiparticle momentum and v is the Fermi velocity (only weakly dependent on the momentum p). The quantum dynamics of quasiparticles can effectively be described by a Dirac-like equation [3]. This spectrum of quasiparticles in graphene has been experimentally verified by observation of specific gate voltage dependencies of Shubnikov–de Haas oscillations, conductivity, and quantum Hall effect [1,2].

Although being different in details, many interesting phenomena in graphene have their analogues in conventional two-dimensional systems. For example, one can observe the quantum Hall effect with a specific structure [1,2] that agrees with theoretical predictions derived from the Dirac equation [4]. Considering effects of disorder one obtains not just the localization but an interesting crossover between the antilocalization and localization behavior [5,6].

At the same time, unique effects specific only for graphene are also possible. One of the most unusual phenomena is the *reflectionless* transmission through a one-dimensional potential barrier of arbitrary strength predicted in Refs. [7,8] and recently coined as the "Klein paradox" [8]. The simplest experimental setup suggested for studying this effect is a graphene-based *n-p* junction that can be made by split-gate technique [1,7,8] (see schematic in Fig. 1). The absence of the backscattering of the massless particles flowing perpendicular to the barrier is related to the chiral nature of them and to a phenomenon of "isospin" conservation [8]. The perfect transmission of the quasiparticles can be explained in a natural way using a standard theory of interband tunneling [9]. Indeed, it has been shown that the transmission probability P through an *n-p* junction is determined by the gap Δ between the electron and hole bands as [9]

$$P \simeq \exp\left[-\frac{\pi\Delta^2}{4\hbar v F}\right], \tag{2}$$

where F is the slope of an x-dependent electrostatic potential in the *n-p* junction. Since the undoped graphene is a gapless material, taking the limit $\Delta \to 0$ leads to the conclusion about the ideal transmission of the quasiparticles flowing perpendicular to the junction.

Such a perfect transmission of quasiparticles through an *n-p* junction might lead to difficulties in confining electrons in future graphene-based electronic devices (like those suggested, e.g., in Ref. [10]). Although in narrow stripes this difficulty can be avoided due to transversal quantization [11], the problem may persist in clean wide 2D samples.

At the same time, the reflectionless penetration is rather sensitive to applying external fields. For example, it is expected [7] that a magnetic field may reduce the transmission or even confine the electrons [12] for certain nonhomogeneous configurations of the field. It is not difficult to imagine that the current flowing through the *n-p* junctions in graphene may not be less sensitive to other external perturbations.

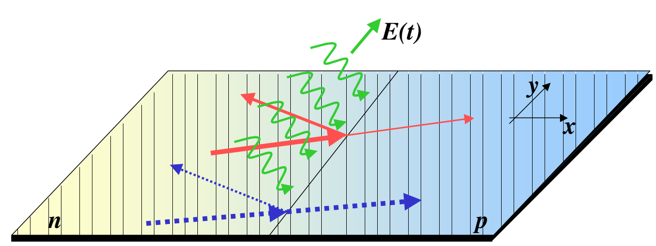

FIG. 1 (color online). An *n-p* junction in graphene. The quasiparticle scattering in the absence (dashed lines) and in the presence (solid lines) of the EF is shown.

In this Letter we predict and analyze a new interesting effect arising in the n-p junctions formed in a wide graphene, irradiated by an external electromagnetic field (EF). The system we consider is represented in Fig. 1. We show that the radiation leads to a very unusual dependence of the dc current I on the voltage V applied across the barrier. At low values of the applied voltage the dependence is linear but the resistivity extracted from it is much larger than the one calculated previously [7]. At certain values of the voltage V the current I through the barrier can even be completely blocked. The entire EF induced current-voltage characteristics (CVC) having the N shape is shown in Fig. 2.

We show that the resonant interaction of the Dirac-like particles with the y component of EF *parallel* to the interface leads to formation of a *nonequilibrium gap* $(2\Delta_R)$ between the electron and hole bands in the quasiparticle spectrum. This gap and specific nonmonotonic dependence of the quasiparticle transmission on the energy ϵ_0 are responsible for the unusual form of CVC in Fig. 2.

Formation of such a dynamic gap is well known for two level systems under radiation [13]. Here we have a continuous spectrum rather than just two levels but the effect persists. Like for the two level systems, one can have a resonance that can be achieved when the frequency ω of EF satisfies the specific condition $\hbar\omega = 2v|\mathbf{p}(x)|$, where $\mathbf{p}(x)$ is the coordinate dependent classical momentum of the quasiparticles. The value of the gap depends strongly on the intensity S and frequency ω of external radiation.

In the presence of EF the two bands time and coordinate dependent Hamiltonian has a form [14]

$$\hat{H}(t) = v\hat{\boldsymbol{\sigma}}\left[\hat{\mathbf{p}} - \frac{e}{c}\mathbf{A}(t)\right] + U(x), \tag{3}$$

where $U(x)$ is the electrostatic potential of the n-p junction, and $\hat{\boldsymbol{\sigma}} = \{\hat{\sigma}_x, \hat{\sigma}_y\}$ are the standard Pauli matrices in the sublattice space. We assume that the potential $U(x)$

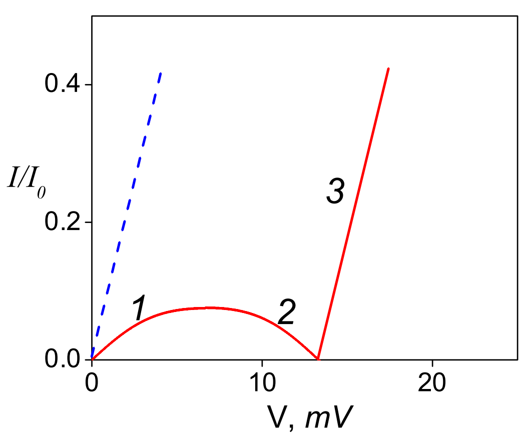

FIG. 2 (color online). The CVC of an n-p junction, in the presence (solid line) and the absence (dashed line) of EF. It is calculated by using Eqs. (13) and (16), and the parameters $\epsilon_0 = 0.02$ eV, $d = 1$ μm, $\hbar\omega = (4/3)\epsilon_0$, and $S = 1$ W/cm^2.

varies sufficiently slowly, such that the scattering between different valleys can be neglected. Just for simplicity and in order to clarify how one comes to the resonant interaction of quasiparticles with the EF, we consider the case of external radiation linearly polarized in the y direction. The electromagnetic wave is characterized by the y component of the vector potential as $A_y = (Ec/\omega)\cos(\omega t)$, and $E = \sqrt{4\pi S/c}$ is the amplitude of the electric field.

Next we reduce the time-dependent problem described by the Hamiltonian (3) to a stationary problem by switching to the rotating frame using the following unitary transformation of the two component Dirac wave functions:

$$\hat{U}_n = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 & 1 \\ -\exp(i\hat{\theta}) & \exp(i\hat{\theta}) \end{pmatrix}\exp\left[i\omega t\left(n - \frac{1}{2}\hat{\sigma}_z\right)\right], \tag{4}$$

where $\hat{\theta} = \tan^{-1}(\hat{p}_y/\hat{p}_x)$. The transformed Hamiltonian $\hat{H}'_{\rm eff} = \hat{U}_n^+ H\hat{U}_n - i\hbar\hat{U}_n^+\dot{\hat{U}}_n$ contains, in general, both static and proportional to $\exp(\pm 2i\omega t)$ parts. However, like for the two level systems [13], in the most interesting regime of the *resonant interaction* between the EF and propagating quasiparticles when

$$\hbar\,\omega \simeq 2v|\mathbf{p}(x)|, \tag{5}$$

only the static part $\hat{H}_{\rm eff}$ is important, and it is written as

$$\hat{H}_{\rm eff} = \begin{pmatrix} \frac{\hbar(2n-1)\omega}{2} + v|\hat{\mathbf{p}}| + U(x) & \frac{eEv}{2\omega} \\ \frac{eEv}{2\omega} & \frac{\hbar(2n+1)\omega}{2} - v|\hat{\mathbf{p}}| + U(x) \end{pmatrix}, \tag{6}$$

where $|\hat{\mathbf{p}}| = \sqrt{\hat{p}_x^2 + \hat{p}_y^2}$. Neglecting the oscillating part of the Hamiltonian $\hat{H}'_{\rm eff}$ corresponds to a rotation wave approximation [13]. A weak nonresonant interaction of quasiparticles with EF neglected here was studied in Ref. [15]. The most important contributions come from almost one-dimensional electron motion, and we assume in our consideration that $p_x \gg p_y$. We also assume that the amplitude of the external microwave radiation is comparatively small, $eEv/\hbar \ll \omega^2$.

Equation (6) shows that the radiation results in the appearance of off-diagonal elements in the operator $\hat{H}_{\rm eff}$. In the absence of the coordinate dependent potential, i.e., $U(x) = 0$, the eigenvalues $\tilde{\epsilon}(p)$ of $\hat{H}_{\rm eff}$ give the sets of bands of quasienergies (the Floquet eigenvalues [13]):

$$\tilde{\epsilon}_{n,\pm}(p) = n\omega \pm \sqrt{(v|\mathbf{p}(x)| - \frac{\hbar\omega}{2})^2 + \Delta_R^2}, \tag{7}$$

where

$$2\Delta_R = (ev/\omega)\sqrt{4\pi S/c} \tag{8}$$

is the EF induced nonequilibrium gap. The n are an integer number $0, \pm 1, \pm 2, \ldots$.

It is well known [13] that, in the presence of periodic time-dependent perturbations, the bands of the Floquet eigenvalues replace the quasiparticle spectrum, Eq. (3). The quantity $\Delta_R/\hbar$ has the same meaning as the famous Rabi frequency for microwave induced quantum coherent oscillations between two energy levels (these energy levels are $v|\mathbf{p}(x)|$ and $-v|\mathbf{p}(x)|$ in our case).

Next we analyze the quasienergy ϵ_0 dependent transmission of quasiparticles $P(\epsilon_0)$ through the potential barrier $U(x)$ formed in the n-p junction. To obtain the analytical solution we use the quasiclassical approximation that can be quite realistic for the n-p junctions created electrostatically. The classical phase trajectories $p(x)$ of the Hamiltonian $\hat{H}_{\text{eff}}$ are determined by the "quasienergy" conservation law as

$$U(x) + \tilde{\epsilon}_{n,\pm}(p) = \epsilon_0. \tag{9}$$

Using Eq. (9) for $n = 0, \pm 1$ we obtain three regimes of the quasiparticle propagation through the barrier. As $\epsilon_0 > \hbar\omega/2$, the resonant condition (5) is satisfied sufficiently close to the junction and, therefore, the interaction with the EF results in the classical reflection of the quasiparticle. This is shown schematically in Fig. 3(a) by a thick solid line. The transmission of the particles through the barrier occurs in the form of the *nonequilibrium interband tunneling* between electron and hole Floquet bands, and therefore, the effect is a particular example of the dynamical tunneling [13]. The quasiparticles tunnel from the electronic $\tilde{\epsilon}_e$ band ($n = 0$) on the left side to the hole $\tilde{\epsilon}_h$ band ($n = 0$) on the right side of the junction. Similarly to the usual case of the interband tunneling [9] the probability $P(\epsilon_0)$ of the quasiparticle transmission is determined by the following process in the "under barrier region": the quasiparticle moves from the left "classical turning point" ($p = \hbar\omega/(2v)$) to the "branch point" ($p = \hbar\omega/(2v) + i\Delta_R/v$) in the complex (x, p) plane, and afterwards to the right "classical turning point."

Further progress can be made by choosing a specific model for the electrostatic potential of voltage biased n-p junction (d is the n-p junction width) [9]

$$U(x) = \begin{cases} eV & x < -d/2 + eV/F \\ F(x + d/2) & -d/2 + eV/F < x < d/2 \\ U = Fd & x > d/2. \end{cases} \tag{10}$$

We obtain for quasiparticles flowing perpendicular to the barrier ($p_y = 0$)

$$P(\epsilon_0) \simeq \exp\left\{i\frac{2}{\hbar}\left[\int_{\epsilon_0-\Delta_R/F}^{\epsilon_0/F} p_+ dx + \int_{\epsilon_0/F}^{\epsilon_0+\Delta_R/F} p_- dx\right]\right\}, \tag{11}$$

where the complex momenta $p_\pm$ are determined by the condition of the quasienergy conservation

$$\epsilon_0 = Fx + \tilde{\epsilon}_e(p_\pm). \tag{12}$$

Calculating the integrals in Eq. (11) we write the transmission probability of the quasiparticles $P(\epsilon_0)$ as [for a particular case as $\epsilon_0 \geq (U - \epsilon_0)$]

$$P(\epsilon_0) \simeq \exp\left[-\frac{\pi\Delta_R^2}{\hbar v F}\right], \qquad (U - \epsilon_0) > \hbar\omega/2, \tag{13}$$

where the gap $2\Delta_R$ should be taken from Eq. (8). Equation (13) shows that the external radiation of the frequency $\omega < 2(U - \epsilon_0)/\hbar$ strongly suppresses the quasiparticle transmission.

In the opposite case of a large frequency ω or small energy ϵ_0, $\omega > 2\epsilon_0/\hbar$, the resonance condition, Eq. (5), cannot be fulfilled, the spectrum remains gapless, and the quasiparticle transmission is not suppressed. In this case the transition occurs from the electronic quasiband with $n = 1$ to the hole quasiband with $n = 0$ [see the gray line in Fig. 3(a)].

There is also a peculiar regime as $(U - \epsilon_0) < \hbar\omega/2 < \epsilon_0$. The interband tunneling for such quasiparticles is *forbidden*, $P \simeq 0$. Indeed, the quasiparticles (electrons) on the left side of the junction starting from the conduction quasiband have to arrive in the forbidden one on the right side of the junction [see Fig. 3(b)].

In experiments, the suppression of the quasiparticle transmission manifests itself as a strong EF induced increase of the resistance of the n-p junction at small voltages $V < [\epsilon_0 - \hbar\omega/2]/e$. The full stationary CVC of the n-p junction in the presence of the EF is determined by the elastic channel; i.e., the energies of quasiparticles on the left and right sides of the junction are equal. Therefore, we write the standard expression for the current I flowing through the ballistic n-p junction as [9]

$$I = \frac{4eL}{(2\pi\hbar)^2}\int d\epsilon dp_y P(\epsilon, p_y)\left[\tanh\frac{\epsilon - eV}{k_BT} - \tanh\frac{\epsilon}{k_BT}\right]. \tag{14}$$

Here, ϵ is the quasienergy of the electrons and $P(\epsilon, p_y)$ is the quasienergy ϵ and p_y dependent transmission of the quasiparticles through the junction. In the case of a wide junction, when the width of the graphene sample L satisfies

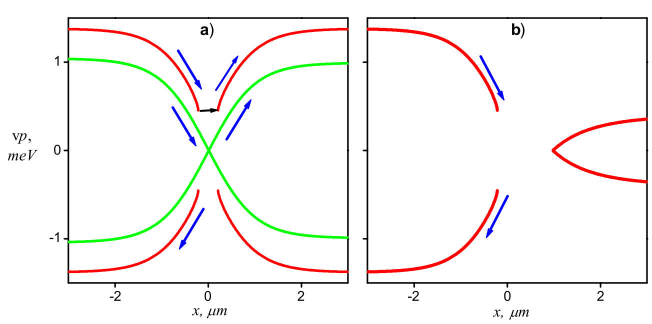

FIG. 3 (color online). The phase trajectories $p(x)$ of the Hamiltonian $\hat{H}_{\text{eff}}$ corresponding to various cases of interband tunneling. (a) Tunneling through the dynamic gap Δ_R (thick solid line, the values of parameters were $U = 2$ meV and $\hbar\omega = 0.8$ meV); the perfect transmission through the point $p = 0$ (gray solid line, $U = 2$ meV and $\hbar\omega = 3$ meV). (b) The zero transmission case ($U = 1.2$ meV and $\hbar\omega = 0.8$ meV). $d = 2$ μm, $\epsilon_0 = 1$ meV, and $\Delta_R = 0.2$ meV were also chosen.

the inequality $L \gg \sqrt{\hbar v d/\epsilon_0}$, the transmission $P(\epsilon, p_y)$ can be written as

$$P(\epsilon, p_y) \simeq P(\epsilon) \exp[-\pi v p_y^2/(\hbar F)], \tag{15}$$

where $P(\epsilon)$ is determined by Eq. (13). Calculating the integral over p_y, and taking the limit of low temperature T, we reduce Eq. (14) for the current I to the form

$$I = I_0 \int_{\epsilon_0}^{\epsilon_0+eV} \frac{d\epsilon}{\epsilon_0} P(\epsilon), \tag{16}$$

where $I_0 = eL\epsilon_0/(\pi\hbar)^2\sqrt{\hbar F/v}$ is the characteristic current of the n-p junction in the absence of the EF [7].

Although the exact shape of the CVC is determined by diverse factors, e.g., by the preexponent in Eq. (13), we argue that the EF with the frequency $\omega < 2\epsilon_0/\hbar$ leads to the N type of the CVC (see Fig. 2). Indeed, as the transport voltage V is less than the characteristic value $V_0 = [\epsilon_0 - \hbar\omega/2]/e$ the quasiparticle current I flows due to the interband tunneling with the probability determined by Eq. (13) [see region 1 in Fig. 2 and the schematic of the process in Fig. 3(a)]. However, in the voltage region $V_0 < V < \min\{2V_0, \epsilon_0\}$ the current I starts to decrease due to the presence of quasiparticles whose propagation is forbidden [see region 2 in Fig. 2 and the schematic in Fig. 3(b)]. The drop of the current becomes especially deep as $2V_0 < \epsilon_0$ and the radiation frequency is in the particular range $\epsilon_0 < \hbar\omega < 2\epsilon_0$. In the voltage region $V > 2V_0$ the current increases with the voltage V because there is a possibility to propagate with the perfect transmission for quasiparticles possessing a small momentum $p < \hbar\omega/(2v)$ [see region 3 in Fig. 2 and the schematic in Fig. 3(b)].

Finally, we address the question of experimental conditions necessary to observe the predicted effects. An n-p junction with the typical width $d \simeq 1\ \mu$m in a graphene sample has to be fabricated. An external radiation containing the component parallel to the junction of a moderate intensity S has to be applied. We emphasize that EF need not be linearly polarized. The EF suppresses the quasiparticle transmission through the n-p junction in the range of the frequencies of EF $\omega \leq \epsilon_0$. This means that for the Fermi energy $\epsilon_0 \simeq 0.02$ eV (this value corresponds to doping levels of a graphene monolayer $n \leq 10^{11}$ cm^{-2} [1,2]), the EF in the far-infrared region with the frequency less than 10^{13} Hz provides a strong decrease of the quasiparticle transmission with the intensity S of the radiation. Choosing an even smaller external frequency $\simeq 2 \times 10^{12}$ Hz one may use the radiation with a moderate intensity $S > 0.4$ W/cm^2 to observe this effect. Notice here that the effect is also reduced at small frequencies $\omega < \epsilon_0\sqrt{v/(\hbar d\epsilon_0)}$ because in this case the transport through the n-p junction is determined by electrons with large values of $p_y > p_x$.

In conclusion, we have demonstrated that radiation of a moderate intensity S having the component parallel to the n-p junction in graphene leads to a pronounced suppression of the quasiparticle transmission through the junction. This effect occurs due to formation of a nonequilibrium dynamic gap between electron and hole bands in the quasiparticle spectrum as the resonant condition (5) is satisfied. Propagation of quasiparticles is possible due to the nonequilibrium interband tunneling. This specific type of the tunneling is determined by the initial energy of electrons ϵ_0 and, as a result, we obtain an N type of CVC. The suppression of the quasiparticle transmission may allow one to control confinement of electrons in diverse structures fabricated in graphene, like, e.g., n-p-n transistors, single electron transistors, quantum dots, etc., by variation of the intensity S and frequency ω of the external radiation. We hope that the predicted effect will find its application to future electronic devices based on graphene.

We would like to thank P. Silvestrov and A. Kadigrobov for useful discussions and acknowledge the financial support by SFB 491 and SFB Transregio 12.

[1] K. S. Novoselov *et al.*, Science **306**, 666 (2004); K. S. Novoselov *et al.*, Nature (London) **438**, 197 (2005).
[2] Y. Zhang, Y.-W. Tan, H. L. Stormer, and P. Kim, Nature (London) **438**, 201 (2005).
[3] G. W. Semenoff, Phys. Rev. Lett. **53**, 2449 (1984).
[4] Y. Zheng and T. Ando, Phys. Rev. B **65**, 245420 (2002); V. P. Gusynin and S. G. Sharapov, Phys. Rev. Lett. **95**, 146801 (2005); F. Guinea, A. H. Castro Neto, and N. M. R. Peres, Phys. Rev. B **73**, 245426 (2006).
[5] H. Suzuura and T. Ando, Phys. Rev. Lett. **89**, 266603 (2002); E. McCann *et al.*, Phys. Rev. Lett. **97**, 146805 (2006).
[6] I. L. Aleiner and K. B. Efetov, Phys. Rev. Lett. **97**, 236801 (2006).
[7] V. V. Cheianov and V. I. Fal'ko, Phys. Rev. B **74**, 041403 (2006).
[8] K. S. Novoselov *et al.*, Nature Phys. **2**, 177 (2006).
[9] E. O. Kane and E. I. Blount, in *Tunneling Phenomena in Solids*, edited by E. Burstein and S. Lundqvist (Plenum, New York, 1969).
[10] B. Trauzettel, D. V. Bulaev, D. Loss, and G. Burkard, Nature Phys. **3**, 192 (2007).
[11] P. G. Silvestrov and K. B. Efetov, Phys. Rev. Lett. **98**, 016802 (2007).
[12] A. De Martino, L. Dell'Anna, and R. Egger, Phys. Rev. Lett. **98**, 066802 (2007).
[13] P. Hänggi, in *Quantum Transport and Dissipation*, edited by T. Dittrich, P. Hänggi, G.-L. Ingold, B. Kramer, G. Schön, and W. Zwerger (Wiley-VCH, Weinheim, 1998).
[14] E. J. Mele, P. Král, and D. Tománek, Phys. Rev. B **61**, 7669 (2000).
[15] B. Trauzettel, Ya. M. Blanter, and A. F. Morpurgo, Phys. Rev. B **75**, 035305 (2007).