Electromagnetic field induced suppression of transport through $n$-$p$ junctions in graphene

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We study quasi-particle transmission through an $n$-$p$ junction in a graphene irradiated by an electromagnetic field (EF). In the absence of EF the electronic spectrum of undoped graphene is gapless, and one may expect the perfect transmission of quasi-particles flowing perpendicular to the junction. We demonstrate that the resonant interaction of propagating quasi-particles with the component of EF parallel to the junction induces a non-equilibrium dynamic gap ($2\Delta_R$) between electron and hole bands in the quasi-particle spectrum of graphene. In this case the strongly suppressed quasi-particle transmission is only possible due to interband tunnelling. The effect may be used for controlling transport properties of diverse structures in graphene, like, e.g., $n$-$p$-$n$ transistors, single electron transistors, quantum dots, etc., by variation of the intensity $S$ and frequency $\omega$ of the external radiation.

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Recent success in fabrication of graphene samples \cite{1,2} has resulted in a stream of publications devoted to study of this interesting material. The unique properties of graphene originate from peculiarities of the electron spectrum. The quasi-particle 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

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

where $p = \{p_x, p_y\}$ is the quasi-particle momentum, $v$ is the Fermi velocity (only weakly dependent on the momentum $p$). The quantum dynamics of quasi-particles can effectively be described by a Dirac- like equation \cite{3}. This spectrum of quasi-particles in graphene has been experimentally verified by observation of specific gate voltage dependencies of Shubnikov-de Haas oscillations, conductivity and quantum Hall effect \cite{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 \cite{1,2} that agrees with theoretical predictions derived from the Dirac equation \cite{4}. Considering effects of disorder one obtains not just the localization but an interesting crossover between the antilocalization and localization behavior \cite{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. \cite{7,8} and recently coined as the “Klein paradox” \cite{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 \cite{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 \cite{5}. The perfect transmission of the quasi-particles can be explained in a natural way using a standard theory of interband tunnelling \cite{6}. Indeed, it has been shown that the transmission probability $P$ through an $n$-$p$ junction is determined by the gap $\Delta$ between the electron and hole bands as \cite{9}

$$P \approx \exp[-\pi\Delta^2/(4\hbar v_F)], \quad (2)$$

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 quasi-particles flowing perpendicular to the junction.

Such a perfect transmission of quasi-particles 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. \cite{10}). Although in narrow stripes this difficulty can be avoided due to transversal quantization \cite{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 \cite{7} that a magnetic field may reduce...
In this Letter we predict and analyze a new interesting regime of the electric field. The externally applied EF is taken into account in the approximation (RWA) [13]. The RWA is valid in 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, \( \epsilon Ev/\hbar \ll \omega^2 \).

The Eq. (3) shows that the radiation results in the appearance of off-diagonal elements in the operator \( \hat{H}_{\text{eff}} \). In the absence of the coordinate dependent potential, i.e. \( U(x) = 0 \), the eigenvalues \( \epsilon(p) \) of \( \hat{H}_{\text{eff}} \) give the sets of bands of quasi-energies (the Floquet eigenvalues [13]):

\[
\epsilon_{n,\pm}(p) = n\omega \pm \sqrt{(\epsilon |p(x)| - \hbar \omega/2)^2 + \Delta_R^2} \tag{7}
\]

A weak non-resonant interaction of quasiparticles with EF neglected here was studied in Ref. [14]. The most interesting regime of the resonant interaction between the EF and propagating quasi-particles when

\[
\hbar \omega \approx 2\epsilon |p(x)| \tag{6}
\]

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

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

where \( \theta = \tan^{-1}(\hat{p}_y/\hat{p}_x) \). The transformed Hamiltonian \( \hat{H} = \hat{H}_0 + \hat{H}_n \) contains, in general, both static and proportional to \( \pm 2\omega t \) parts. However, like for the two level systems [13], only the static part \( \hat{H}_{\text{eff}} \) is important near the resonance, and can be written as

\[
\hat{H}_{\text{eff}} = \left( \frac{\hbar (2n+1)\omega}{2} + \epsilon |p| + U(x) - \frac{\epsilon Ev}{2\omega} \right) - \frac{\hbar (2n+1)\omega}{2} - \frac{\epsilon Ev}{2\omega} \right),
\]

where \( |p| = \sqrt{p_x^2 + p_y^2} \). Neglecting the oscillating part of the Hamiltonian \( \hat{H} \), it corresponds to a rotation wave approximation (RWA) [13]. The RWA is valid in the most interesting regime of the resonant interaction between the EF and propagating quasi-particles when

\[
\hbar \omega \approx 2\epsilon |p(x)| \tag{6}
\]

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The Eq. (3) shows that the radiation results in the appearance of off-diagonal elements in the operator \( \hat{H}_{\text{eff}} \). In the absence of the coordinate dependent potential, i.e. \( U(x) = 0 \), the eigenvalues \( \epsilon(p) \) of \( \hat{H}_{\text{eff}} \) give the sets of bands of quasi-energies (the Floquet eigenvalues [13]):

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\epsilon_{n,\pm}(p) = n\omega \pm \sqrt{(\epsilon |p(x)| - \hbar \omega/2)^2 + \Delta_R^2} \tag{7}
\]
where
\[ 2\Delta_R = (ev/\omega)\sqrt{4\pi S/c} \] (8)
is the EF induced non-equilibrium 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 quasi-particle spectrum, Eq. (1).

The quantity \( \Delta_R/h \) has the same meaning as the famous Rabi frequency for microwave induced quantum coherent oscillations between two energy levels (these energy levels are \( v|p(x)| \) and \( -v|p(x)| \) in our case).

Next we analyze the quasi-energy \( \epsilon_0 \) dependent transmission of quasi-particles \( P(\epsilon_0) \) through the potential barrier \( U(x) \) formed in the \( n-p \) junction. To obtain the analytical solution we use the quasi-classical approximation that can be quite realistic for the \( n-p \) junctions created electrostatically. The classical phase trajectories \( p(x) \) of the Hamiltonian \( H_{eff} \) are determined by the conservation of the sum of the potential energy \( U(x) \) and the quasi-energy \( \epsilon_{n,\pm}(p) \) as
\[ U(x) + \epsilon_{n,\pm}(p) = \epsilon_0 . \] (9)

Using Eq. (9) for \( n = 0, \pm 1 \) we obtain three regimes of the quasi-particle propagation through the barrier depending on the energy of electrons \( \epsilon_0 \) and frequency \( \omega \).

As \( \epsilon_0 > \hbar \omega/2 \), the momentum \( p_x \) decreases as the quasi-particle approaches the barrier, the resonant condition is satisfied sufficiently close to the junction and, therefore, the interaction with the EF results in the classical reflection of the quasi-particle. This is shown schematically in Fig. 3a by a thick solid line. The transmission of the particles through the barrier occurs in the form of the non-equilibrium interband tunnelling between electron and hole Floquet bands, and therefore, the effect is a particular example of the dynamical tunnelling \[13\]. 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 tunnelling \[10\] the probability \( P(\epsilon_0) \) of the quasi-particle transmission is determined by the following process in the “under barrier region” : the quasi-particle moves from the left “classical turning point” \( (p = \hbar \omega/(2v)) \) to “the branch point” \( (p = \hbar \omega/(2v) + i2\Delta_R/v) \) in the complex \( (x, p) \) plane, and afterwards to the right “classical turning point”.

A 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) \[9\]
\[ U(x) = \begin{cases} 
  eV, & x < -d/2 + eV/F \\
  F(x + d/2), & -d/2 + eV/F < x < d/2 \\
  U = Fd, & x > d/2 
\end{cases} \] (10)

We obtain for quasi-particles flowing perpendicular to the barrier \( (p_y = 0) \):
\[ P(\epsilon_0) \simeq \exp \left\{ \frac{2}{\hbar} \left[ \int_{\omega-\epsilon_0}^{\omega+\epsilon_0} p_+ dx + \int_{\omega-\epsilon_0}^{\omega+\epsilon_0} p_- dx \right] \right\} , \] (11)

where the complex momenta \( p_\pm \) are determined by the condition of the quasi-energy conservation
\[ \epsilon_0 = \hbar \omega + \tilde{\epsilon}_e(p_\pm) . \] (12)

Calculating the integrals in Eq. (11) we write the transmission probability of the quasiparticles \( P(\epsilon_0) \) as (for a particular case as \( \epsilon_0 \geq (U - \epsilon_0) \))
\[ P(\epsilon_0) \simeq \exp \left[ -\frac{\pi \Delta_R^2}{h^2 \omega^2} \right] (U - \epsilon_0) > \hbar \omega/2 , \] (13)

where the gap \( 2\Delta_R \) should be taken from Eq. (8). The Eq. (13) shows that the external radiation of the frequency \( \omega < 2(U - \epsilon_0)/h \) strongly suppresses the transmission of the quasiparticles flowing perpendicular to the junction.

In the opposite case of a large frequency \( \omega \) or small energy \( \epsilon_0 \), \( \omega > 2\epsilon_0/h \), the resonance condition, Eq. (6), cannot be fulfilled, the spectrum remains gapless and the quasiparticle transmission is not suppressed. In this case the transition occurs from the electronic quasi-band with \( n = 1 \) to the hole quasi-band with \( n = 0 \) (see the gray line in Fig. 3a).

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 quasi-particles (electrons) on the left side of the junction starting from the conduction quasi-band have to arrive in the forbidden one on the right side of the junction (see Fig. 3b).

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 transport voltages \( V < \left| \epsilon_0 - \hbar \omega/2 \right|/e \). 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 can write the standard expression for the current \( I \) flowing through the ballistic \( n-p \) junction as \[9\]
\[ I = \frac{4eL}{(2\pi\hbar)^2} \int d\epsilon dp_y P(\epsilon, p_y) \left[ \tanh \frac{\epsilon - eV}{k_B T} - \tanh \frac{-\epsilon}{k_B T} \right] . \] (14)

Here, \( \epsilon \) is the quasi-energy of the electrons and \( P(\epsilon, p_y) \) is the quasi-energy \( \epsilon \) and \( p_y \) dependent transmission of the quasiparticles through the junction. The coefficient 4 in Eq. (14) is due to the spin and valleys degeneracy of the quasi-particle spectrum in graphene. In the case of a wide junction, when the width of the graphene sample \( L \)
the voltage region $V > \epsilon_0$. However, in the voltage region due to the presence of quasi-particles whose propagation is forbidden (see, the region 2 on the CVC in Fig. 2), the transmission $P(\epsilon, p_y)$ can be written as

$$P(\epsilon, p_y) \simeq P(\epsilon) \exp\left[-\pi v p_y^2/(\hbar F)\right],$$

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

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

where $I_0 = eL/\pi v \hbar \sqrt{F/\epsilon_0}$ is the characteristic current flowing through the $n$-p junction in the absence of the EF.

Although the exact shape of the CVC is determined by diverse factors, e.g., by the pre-exponent in Eq. (13), and therefore, by the particular form of the electrostatic potential, temperature etc., we argue that the EF with the frequency $\omega < 2\hbar \epsilon_0/\hbar$ 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 tunnelling with the probability determined by Eq. (13) (see, the region 1 on the CVC in Fig. 2 and corresponding schematic of the process in Fig. 3a). However, in the voltage region $V_0 < V < \min\{2V_0, \epsilon_0\}$ the current $I$ starts to decrease due to the presence of quasi-particles whose propagation is forbidden (see, the region 2 on the CVC in Fig. 2 and the schematic in Fig. 3b). 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 quasi-particles possessing a small momentum $p < \hbar \omega/(2v)$ (see the region 3 on the CVC and the schematic in Fig. 3b).

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 quasi-particles 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 $\epsilon_0 \simeq 0.02eV$ (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 quasi-particle transmission with the intensity $S$ of the radiation. Choosing an even smaller external frequency $\omega \simeq 2 \cdot 10^{12}Hz$ and the width of an $n$-p junction $d = 1\mu m$ 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 $\omega < \epsilon_0 \sqrt{v/(\hbar \epsilon_0)}$ 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 quasi-particle transmission through the junction. This effect occurs due to formation of a non-equilibrium dynamic gap between electron and hole bands in the quasi-particle spectrum as the resonant condition $\omega = 2v|\mathbf{p}|$ is satisfied. The value of the gap can be controlled by variation of the intensity $S$ of an external radiation. Propagation of quasiparticles is possible due to the non-equilibrium interband tunnelling. This specific type of the tunnelling is determined by the initial energy of electrons $\epsilon_0$ and the frequency of EF $\omega$ and, as a result we obtain an $N$-type of CVC. The suppression of the quasiparticle transmission may allow 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 $\omega$ of the external radiation. We hope that the predicted effect will find its application to future electronic devices based on graphene.

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