

Magnetophotoelectric effect in graphene via tailored potential landscapes

Joris Josiek^{1,2}, Friedemann Queisser^{1,2}, Stephan Winnerl¹, and Ralf Schützhold^{1,2}

¹*Helmholtz-Zentrum Dresden-Rossendorf, Bautzner Landstraße 400, 01328, Dresden, Germany*

²*Institut für Theoretische Physik, Technische Universität Dresden, 01062, Dresden, Germany*



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We study the propagation of charge carriers in planar graphene under the combined influence of a constant transversal magnetic field B and an in-plane varying electric potential $\phi(x)$. By suitably designing the potential landscape $\phi(x)$, we may effectively steer charge carriers generated by photoexcitation, for example, in order to achieve an efficient charge separation. In comparison to previously considered scenarios for the magnetophotoelectric effect in graphene involving folded graphene or graphene edges, for example, this scheme offers advantages in terms of robustness, scalability, and control. Thus, our findings may pave the way for transport schemes (“graphene electronics”) or photoelectric/photovoltaic applications.

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I. INTRODUCTION

The special properties of graphene, such as the high charge-carrier mobility and its long mean free path, suggest utilizing it for photoelectric or photovoltaic applications [1–20]. In contrast to semiconductors, for example, graphene has no intrinsic band gap. On the one hand, this can be advantageous since restrictions such as the Shockley-Queisser limit do not necessarily apply [21]. On the other hand, it requires alternative schemes for charge separation (see also Refs. [22,23]). Obviously, one has to break the $\mathcal{C}$ (charge), $\mathcal{P}$ (parity), and $\mathcal{T}$ (time reversal) symmetries in order to achieve charge separation. The simplest way to break all these symmetries would be an electric field (e.g., via a bias voltage), but this alone would not facilitate photovoltaic applications without further ingredients (such as a band gap). Another way to break these symmetries is to apply an external magnetic field, but this is also not enough. In order to turn the usual circular cyclotron orbits in a magnetic field into a directed motion of the charge carriers, one can use, for example, geometric constraints such as folded graphene [23] or a graphene edge [22,24,25].

However, such geometric constraints go along with experimental and technological challenges. Graphene edges are typically quite vulnerable to imperfections, and it is very hard to scale up these schemes based on graphene folds or edges to cover large enough areas.

In order to avoid these difficulties, we consider an alternative scheme based on an additional electrostatic potential $\phi(x)$. This could be induced externally by electrodes exhibiting gate voltages or internally by effectively doping the graphene sheet, which shifts the local chemical potentials (see also Refs. [26–34]). The idea is to use the interplay between the magnetic field B and the electrostatic potential $\phi(x)$ in order to systematically manipulate the nonequilibrium dynamics of the charge carriers generated via photoexcitation. To reach this goal, we have to understand the structure of the electron and hole solutions in dependence of B and $\phi(x)$, e.g., its symmetries.

In the following, we tackle this problem by studying the dynamics of charge carriers in graphene under the combined influence of a constant transversal magnetic field B and an in-plane varying electric potential $\phi(x)$. To this end, we start from the effective Dirac equation, i.e., a quantum treatment, in Sec. II. In order to visualize and understand the behavior of the solutions, we discuss the associated semiclassical trajectories in Sec. III. These results allow us to study the $\mathcal{C}$, $\mathcal{P}$, and $\mathcal{T}$ symmetries mentioned above in dependence on the properties of the electrostatic potential $\phi(x)$ in Sec. IV. We find that by suitably tailoring the potential $\phi(x)$, one should be able to achieve the charge separation required for a magnetophotoelectric effect (see Sec. V). In order to discuss the applicability of our scheme, we consider some exemplary experimental parameters in Sec. VI.

As already anticipated above, the scheme presented here offers potential advantages in comparison to previous setups based on geometric constraints such as graphene folds or edges. First, it can be realized based on an undisturbed planar graphene sheet and is thus less prone

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to defects and imperfections than, for example, a graphene edge. Second, it is much easier to scale up this scheme to larger areas. Third, it offers more options for manipulation and control. Apart from varying the applied magnetic field B , one can adjust the electrostatic potential landscape $\phi(x)$ by changing the gate voltages or switching on or off electrodes.

II. DIRAC EQUATION

Even though the main idea of the underlying mechanism can already be understood in terms of a semiclassical approach (see also Sec. III), let us start with a quantum description. This allows us to estimate the range of applicability of this semiclassical approach as well as deviations from it, such as tunneling through the potential barriers.

On length scales far above the lattice spacing of approximately 0.25 nm and energies well below the hopping energy approximately equal to 2.8 eV, we may describe electrons and holes by an effective Dirac equation in $2 + 1$ dimensions ($\hbar = 1$)

$$\gamma^\mu (\partial_\mu + i q A_\mu) \Psi = 0, \quad (1)$$

with $x^\mu = [v_F t, x, y]$, where $v_F \approx 10^6$ m/s is the Fermi velocity [35]. The Dirac matrices $\gamma^\mu = [\sigma^z, i\sigma^y, -i\sigma^x]$ acting on the spinor $\Psi = [\psi_1, \psi_2]$ are related to the Pauli matrices $\sigma^{x,y,z}$. In this representation, the Dirac α and β matrices are given by $\alpha^x = \sigma^x$, $\alpha^y = \sigma^y$, and $\beta = \sigma^z$. The $2 + 1$ vector potential $A_\mu = [\phi(x), 0, Bx]$ in the Landau gauge generates the electric field $E(x)$ in the x direction and the magnetic field B in the z direction.

In view of the translation symmetry in t and y , we can make the usual separation ansatz for the modes

$$\Psi(t, x, y) = \exp \{-i\omega t + iky\} \Psi^{\omega,k}(x), \quad (2)$$

arriving at the two coupled equations

$$\begin{aligned} v_F [\partial_x + k + qBx] \psi_2^{\omega,k}(x) &= i[\omega - q\phi(x)] \psi_1^{\omega,k}(x), \\ v_F [\partial_x - k - qBx] \psi_1^{\omega,k}(x) &= i[\omega - q\phi(x)] \psi_2^{\omega,k}(x). \end{aligned} \quad (3)$$

Without loss of generality, we choose $\psi_1^{\omega,k}(x)$ to be real while $\psi_2^{\omega,k}(x)$ is imaginary.

A. Dispersion relation

In order to obtain a first insight into the structure of the solutions to Eq. (3), we employ WKB analysis (see the Appendix) and arrive at the dispersion relation

$$p_x^2(x) + [k + qBx]^2 = \frac{[\omega - q\phi(x)]^2}{v_F^2}, \quad (4)$$

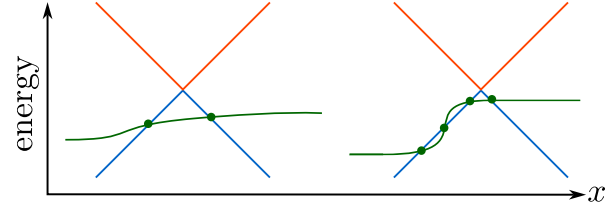


FIG. 1. Sketch of solutions to the dispersion relation (4) for the subthreshold case (left) and the superthreshold case (right). The green curve denotes the potential $q\phi(x)$, while the red and blue straight lines correspond to $\omega \pm v_F[k + qBx]$. In the subthreshold case (left), the classically allowed solution lies in between the two intersections (green dots) with the two straight lines, which is referred to as type I. In the superthreshold case (right), there is an additional classically allowed solution lying in between two intersections with the *same* straight line, which is referred to as type II.

where $p_x(x)$ denotes the momentum in the x direction. The classical turning points where $p_x = 0$ can thus be obtained from the intersections between the potential curve $\phi(x)$ and straight lines with slope $\pm v_F B$, shifted by k and ω , i.e., $q\phi(x) = \omega \pm v_F[k + qBx]$ (see Fig. 1).

B. Subthreshold fields

Assuming that the potential $\phi(x)$ approaches constant values asymptotically $\phi(x \rightarrow \pm\infty) = \text{const}$, the possible number of intersections depends on the maximum slope of $\phi(x)$, i.e., the electric field $E(x)$. If this field is subthreshold $E(x) < v_F B$ everywhere, we always find two intersections—one for each straight line (see Fig. 1, left). In the following, we refer to these solutions as type I.

In the limit of a vanishing electric field $\phi = \text{const}$, this case just corresponds to the usual circular (cyclotron) motion [see Fig. 2(a)]. The distance between the two intersections is then twice the cyclotron radius

$$r_{\text{cyc}} = \frac{\omega - q\phi}{qBv_F}. \quad (5)$$

Adding a small and constant electric field induces a drift velocity in the orthogonal direction $\mathbf{v}_{\text{drift}} = \mathbf{E} \times \mathbf{B} / B^2$. Exploiting the quasi-relativistic form of Eq. (1), this can be understood via applying an effective Lorentz boost with $\mathbf{v}_{\text{drift}}$ after which the electric field in the comoving frame vanishes. However, since this drift velocity is the same for particles and holes, it cannot be used to achieve charge separation directly.

C. Superthreshold fields

Let us now consider an electric field that exceeds the threshold $E(x) > v_F B$ in some region. In the limiting case of constant fields with $E > v_F B$, an effective Lorentz boost could transform the magnetic field away.

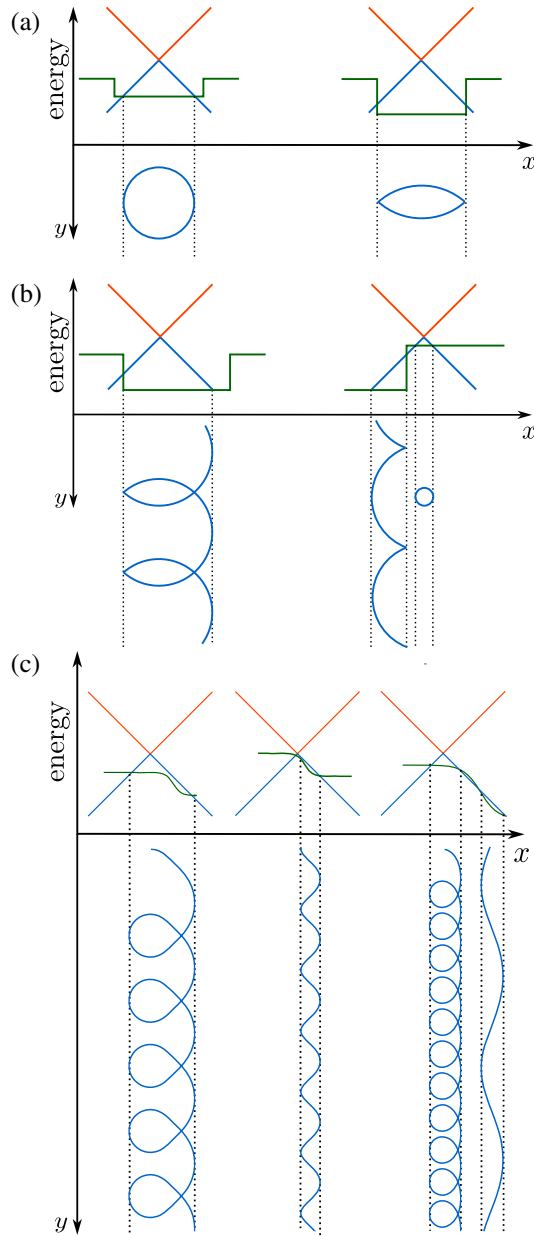


FIG. 2. Sketch of various solutions to the dispersion relation (4) together with the associated semiclassical trajectories. (a) The cyclotron motion driven solely by the magnetic field is sketched in the left panel, while the right panel shows a localized trajectory with turning points set by the potential $\phi(x)$. (b) A combination of these two mechanisms for turning is depicted in the left panel. So far, all of these solutions were of type I. The right panel shows the coexistence of a type-II solution, where $v_y \neq 0$, with a localized cyclotron motion of type I. In contrast to the steplike behavior of the potential in panels (a) and (b), analogous solutions (obtained numerically) for smooth potentials are shown in panel (c). The left panel displays a type-I solution, where both $\dot{x}(t)$ and $\dot{y}(t)$ change sign. In the middle panel, the type-II solution exhibits turning points only in $x(t)$. The right panel illustrates the coexistence of type-I and type-II solutions.

Note that the same condition $E > v_F B$ is required for observing the graphene analog of the Sauter-Schwinger effect, i.e., electron-positron pair creation of the vacuum in quantum electrodynamics (QED) induced by a strong (and slowly varying) electric field [36,37]. However, this field strength $E > v_F B$ should not be confused with the analog of the Schwinger critical field in QED

$$E_{\text{crit}} = \frac{m_e^2 c^3}{\hbar q} \approx 1.3 \times 10^{18} \frac{\text{V}}{\text{m}}, \quad (6)$$

where m_e is the electron mass and c the speed of light. Since graphene does not have a band gap, there is no direct analog to this quantity (6). However, as one may infer from Eq. (4), for example, a transversal momentum $k_\perp$ could act like an effective gap $v_F k_\perp$. If we assume such a gap in the eV regime, the graphene analog of the above critical field (6) would be of order 10^9 V/m. In contrast, even for a magnetic field of 1 T, the field strength required for reaching $E > v_F B$ is much lower $E = \mathcal{O}(10^6 \text{ V/m})$. Generating such a field—e.g., via a potential difference $\Delta\phi = \mathcal{O}(\text{V})$ over a submicrometer scale—is much easier than reaching 10^9 V/m. Note that a sharp potential step such as $\phi(x) = \Delta\phi \Theta(x)$ would correspond to the limit of infinite field strengths and could allow us to observe the analog of Klein tunneling in graphene, see, e.g., Refs. [38,39].

For superthreshold fields $E(x) > v_F B$, there may be additional intersections that now lie on the *same* straight line, as in Fig. 1 right. These additional solutions will be referred to as type II.

Since both intersections lie on the same straight line (e.g., the line with positive slope in Fig. 1 right), the propagation velocity in the y direction, which is determined by the mechanical momentum $k + qBx$, does not change sign in between these turning points—which means that the motion in the y direction is not reversed (in contrast to the circular orbits discussed above). As an intuitive picture, one turning point is caused by the bending due to the magnetic field, while the other turning point is caused by a reflection at the electrostatic potential barrier.

If the electric field is above the threshold $E > v_F B$ but still far below the critical value of 10^9 V/m discussed above, the tunneling probability through the electrostatic potential barrier is exponentially suppressed unless the incoming energy becomes close to the top of the barrier.

As one can already infer from Fig. 1, one can have a type-I solution without a type-II solution, but not the other way around (though they may lie on different branches). In general, one can have a coexistence of the different trajectories, as in Fig. 2, bottom right.

D. Quantum solutions

On the classical level, the case of two intersections with $p_x^2 > 0$ in between corresponds to a bound state in the x direction that could, however, display a nonzero average velocity in the y direction. If the size L of the classically allowed region where $p_x^2 > 0$ is large enough, i.e., $p_x^2 L^2 \geq \mathcal{O}(\hbar^2)$, these states will also exist on the quantum level, i.e., as solutions to Eq. (3). Neglecting the exponentially suppressed tunneling through the potential barrier (as discussed above), these solutions would be analogous to a discrete set of bound states in the x direction. The associated stationary states will be localized in the x direction (with exponential tails in the classically forbidden region $p_x^2 < 0$) and plane waves in the y direction, characterized by a discrete set of frequencies $\omega(k)$. In principle, given the quantization of frequencies ω (for a fixed value of k), one could obtain resonance effects, such as resonant enhancement of photon absorption. However, for characteristic parameters (see Sec. VI below), there are typically many discrete levels lying quite close together such that we may approximate them as a semiclassical continuum. The motion in the y direction can either be obtained from the group velocity $d\omega/dk$ or the current, where we can apply many of the arguments already discussed in Ref. [23] after incorporating the additional potential $\phi(x)$.

To be more explicit, let us consider the current $j_{\omega,k}^\mu$ associated to an eigensolution $\Psi^{\omega,k}$, which reads

$$j_{\omega,k}^\mu = v_F \bar{\Psi}^{\omega,k} \gamma^\mu \Psi^{\omega,k}. \quad (7)$$

The density $\rho = |\psi_1^{\omega,k}|^2 + |\psi_2^{\omega,k}|^2$ is recovered for $\mu = 0$. As expected for stationary solutions that are localized in the x direction, the $\mu = 1$ contribution, i.e., the current $j_{\omega,k}^x = v_F \Psi^{\omega,k} \sigma^x \Psi^{\omega,k} = v_F (\psi_1^{\omega,k})^* \psi_2^{\omega,k} + \text{H.c.} = 0$ vanishes identically since $\psi_1^{\omega,k}$ is real, while $\psi_2^{\omega,k}$ is imaginary.

The remaining current in the y direction,

$$j_{\omega,k}^y = v_F \Psi^{\omega,k} \sigma^y \Psi^{\omega,k} = i v_F (\psi_2^{\omega,k})^* \psi_1^{\omega,k} + \text{H.c.}, \quad (8)$$

can be simplified to $j_{\omega,k}^y = -2i v_F \psi_1^{\omega,k} \psi_2^{\omega,k}$. Assuming that ω lies outside the range of $q\phi(x)$, as in Fig. 1, we may derive the total current in the y direction by inserting Eq. (3) and integrating over x , which gives the two equivalent expressions

$$\begin{aligned} J_{\omega,k}^y &= \int dx j_{\omega,k}^y \\ &= 2v_F^2 \int dx \left[\frac{k + qBx}{\omega - q\phi} + \frac{q\phi'}{2[\omega - q\phi]^2} \right] |\psi_1^{\omega,k}|^2 \\ &= 2v_F^2 \int dx \left[\frac{k + qBx}{\omega - q\phi} - \frac{q\phi'}{2[\omega - q\phi]^2} \right] |\psi_2^{\omega,k}|^2. \end{aligned} \quad (9)$$

The second term $\propto q\phi'$ stems from the commutator of ∂_x and $q\phi(x)$, i.e., it is an effect of the full quantum treatment and can be neglected in the semiclassical limit $p_x^2 L^2 \gg \hbar^2$. Furthermore, it comes with opposite signs depending on the representation (in terms of $\psi_1^{\omega,k}$ or $\psi_2^{\omega,k}$). Thus, the direction (i.e., sign) of the current $J_{\omega,k}^y$ is determined by the first term, which is the same in both representations and coincides with the results of the semiclassical approach.

E. Charge separation

To infer the direction of the current $J_{\omega,k}^y$, we have to study the relative sign between $k + qBx$ and $\omega - q\phi(x)$. For type-I solutions, the sign of $k + qBx$ changes in between the turning points such that the integral could be positive or negative or even zero (which is the case for circular orbits in a pure magnetic field). For type-II solutions, on the other hand, $k + qBx$ does not change sign in the classically allowed region.

As a result, the associated bound states in the x direction move with a finite velocity in the y direction (assuming that one can neglect the exponentially suppressed contributions of the evanescent tails). Hence, these modes are very relevant for achieving charge separation. Even if ω lies inside the range of $q\phi(x)$, we may draw the same conclusion on the semiclassical level because the point where $\omega = q\phi(x)$ is also outside the classically allowed region.

The above conclusion based on the current can also be applied to the group velocity $d\omega/dk$. To show this, let us write Eq. (3) in terms of the Dirac Hamiltonian $\hat{H}_k$ per mode k in the form $\hat{H}_k |\Psi^{\omega,k}\rangle = \omega |\Psi^{\omega,k}\rangle$ where

$$\hat{H}_k = \begin{pmatrix} q\phi(x) & -i v_F [\partial_x + k + qBx] \\ -i v_F [\partial_x - k - qBx] & q\phi(x) \end{pmatrix}. \quad (10)$$

Taking the k derivative $d\hat{H}_k/dk = v_F \sigma^y$, we see that the current J^y can be written as $\langle \Psi^{\omega,k} | d\hat{H}_k/dk | \Psi^{\omega,k} \rangle$. Next, taking the k derivative of $\hat{H}_k |\Psi^{\omega,k}\rangle = \omega |\Psi^{\omega,k}\rangle$, we find that the current J^y coincides with the group velocity $d\omega/dk$ for normalized solutions $|\Psi^{\omega,k}\rangle$.

The triangle inequality $2|ab| \leq |a|^2 + |b|^2$ implies that $|j^y| \leq v_F \rho$, i.e., the group velocity $d\omega/dk$ cannot exceed the Fermi velocity v_F , as expected.

III. SEMICLASSICAL TRAJECTORIES

In order to visualize the type-I and type-II solutions discussed above, let us study the associated semiclassical trajectories $\mathbf{r}(t) = [x(t), y(t)]$ (see the Appendix). This allows us to understand how the motion of the charge carriers is affected by the interplay between the magnetic field B and the potential landscape $\phi(x)$. We may then use

this insight to design potential landscapes $\phi(x)$ that facilitate a strong photoelectric effect. As explained above, for suitable parameters, the results of this semiclassical analysis can also be applied (at least qualitatively) to the full quantum treatment.

Since the effective Dirac Eq. (1) does not contain a mass term, we are in the ultrarelativistic limit where the velocity is given by

$$\dot{\mathbf{r}} = v_F \frac{\mathbf{p}}{|\mathbf{p}|}, \quad (11)$$

with the momentum $\mathbf{p}$ satisfying the equation of motion

$$\dot{\mathbf{p}} = q(\mathbf{E} + \dot{\mathbf{r}} \times \mathbf{B}) \quad (12)$$

(see the Appendix). Solving this coupled set of equations numerically yields the trajectories $\mathbf{r}(t)$, where some example cases are depicted in Fig. 2(c).

The group velocity $d\omega/dk$ or the current J^y of the stationary solutions $\Psi^{\omega,k}$ is then related to the average velocity v_y on the classical level. For the type-II solutions, the velocity $v_y(t)$ displays some time dependence but never changes its sign—such that the average velocity v_y is also nonzero. For the type-I solutions, on the other hand, the velocity $v_y(t)$ does change its sign during each cycle. However, we found that most of them also yield a non-vanishing average velocity v_y . Only in quite special cases did it average out to zero.

To understand the origin of this average velocity v_y , let us consider the simplified case of rectangular potential barriers (see Figs. 3 and 4). We see that the average velocity v_y only vanishes for the symmetric scenarios where the contributions from the two branches $\omega \pm v_F[k + qBx]$ exactly cancel each other. Note that, strictly speaking, rectangular potential barriers would imply diverging electric fields such that special care is required when applying some of the above methods, such as the WKB approximation. Instead, one could start with the known solutions for flat potentials $\partial_x \phi = 0$ and connect them via suitable boundary conditions. More realistically, however, these rectangular potential barriers are just approximations for smoothly varying landscapes $\phi(x)$.

IV. SYMMETRIES

Now let us investigate the solutions discussed above in light of the $\mathcal{C}$, $\mathcal{P}$, and $\mathcal{T}$ symmetries already mentioned in the Introduction. Applied separately, each of these transformations would reverse the electric and/or magnetic field and thus change the setup under consideration. Thus, we shall try to find a combined symmetry that keeps the electric and magnetic fields fixed. To this end, we have to specify how the potential $\phi(x)$ transforms under reflection $x \rightarrow -x$.

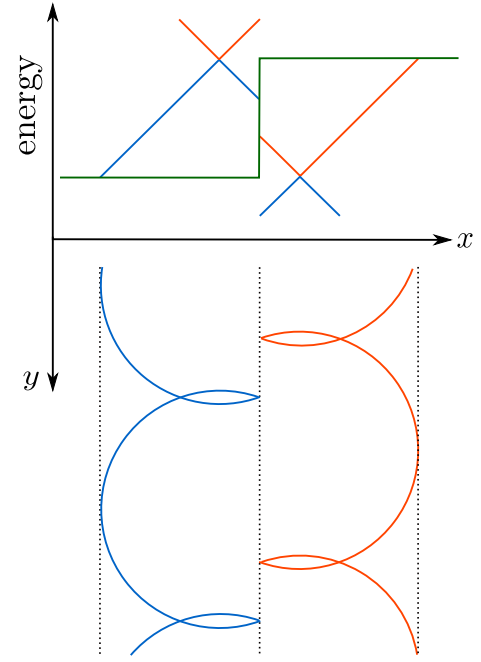


FIG. 3. Example of an odd potential $\phi(x)$ that permits the existence of paired electron and hole solutions that propagate in the same y direction.

A. Odd potential

First, we assume that the potential is an odd function $\phi(-x) = -\phi(x)$, such as a potential step (see Fig. 3). Then we find that Eq. (3) remains invariant under the combined reflection $x \rightarrow -x$ as well as $k \rightarrow -k$ and $\omega \rightarrow -\omega$. Since

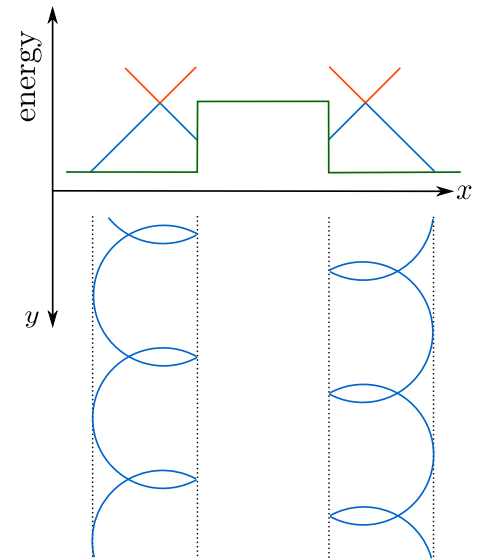


FIG. 4. Example of an even potential $\phi(x)$ that permits the existence of paired solutions with the same frequency (either both electron solutions or both hole solutions) that propagate in opposite directions.

both k and ω change sign, the group velocity (along the potential, i.e., in the y direction) stays the same.

For example, if we have an electron solution with $\omega > 0$ (assuming a vanishing chemical potential $\mu = 0$) that is localized on the right-hand side of the potential step and moves in the y direction, the above symmetry implies the existence of a hole solution $\omega < 0$ that is localized on the left-hand side of the potential step and moves in the same y direction (see Fig. 3).

B. Even potential

Now let us assume an even function $\phi(-x) = \phi(x)$ such as a potential barrier (see Fig. 4). In this case, Eq. (3) remains invariant if we apply the reflection $x \rightarrow -x$ as well as $k \rightarrow -k$ (but now ω remains fixed) and simultaneously transform the spinor wave function via $\psi_2^{\omega,k} \rightarrow -\psi_2^{\omega,k}$, i.e., $\Psi^{\omega,k} \rightarrow \sigma^z \Psi^{\omega,k} = (\Psi^{\omega,k})^*$.

In this case, a solution that is localized on the right-hand side of the potential barrier and moves in the y direction, for example, is transformed into another solution with the same frequency that is localized on the left-hand side of the potential barrier and moves in the opposite y direction (see Fig. 4).

At the same time, a solution with the opposite frequency can move on a qualitatively different trajectory, e.g., it could experience the potential barrier as a potential dip, which can be overcome easily.

V. MAGNETOPHOTOELECTRIC EFFECT

Now we have all the ingredients required for designing a potential landscape that facilitates an effective charge separation. First, we found a definite propagation direction (at least for all type-II solutions, but also for most of the type-I cases) orthogonal to the electric and magnetic fields (i.e., in the y direction). Second, the trajectories are different for electrons and holes. Apart from the opposite orientation of their (quasi) circular orbits in the magnetic field, a negative potential barrier $\phi(x) < 0$ acts as a reflector for electrons (as long as it is high and wide enough) while holes can traverse it—while it is the other way around for a positive potential barrier $\phi(x) > 0$.

An example of such a potential landscape $\phi(x)$ is sketched in Fig. 5. By comblike structures of alternating positive and negative potential barriers, we can steer the charge carriers. The majority of the electron trajectories (with large enough cyclotron radii) encircle the negative potential barriers in a counterclockwise direction, while the hole trajectories encircle the positive potential barriers in a clockwise direction.

To collect the charge carriers, one could envision extending the positive potential barriers slightly upward relative to the negative potential barriers (i.e., breaking the translational invariance in the y direction at some point) and attaching electric contacts to the graphene that absorb

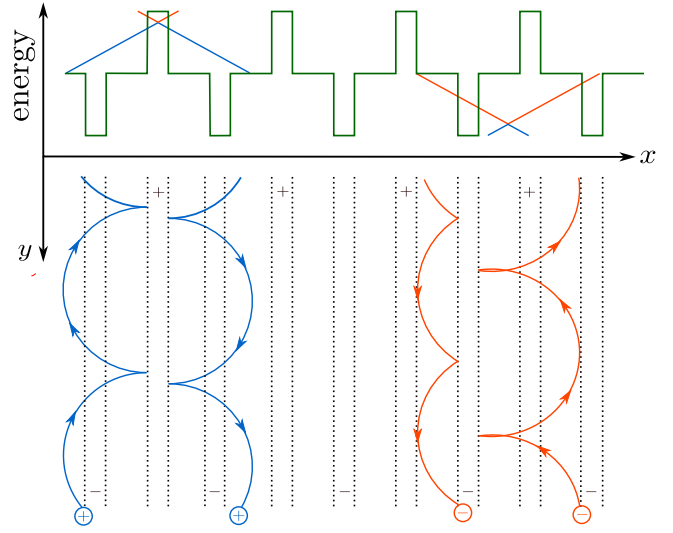


FIG. 5. Arranging positive and negative potential barriers in a double comblike structure allows us to separate the charge carriers. Negative potential barriers reflect electrons but can be traversed by holes. Conversely, holes are reflected by positive potential barriers while being able to propagate through negative potentials.

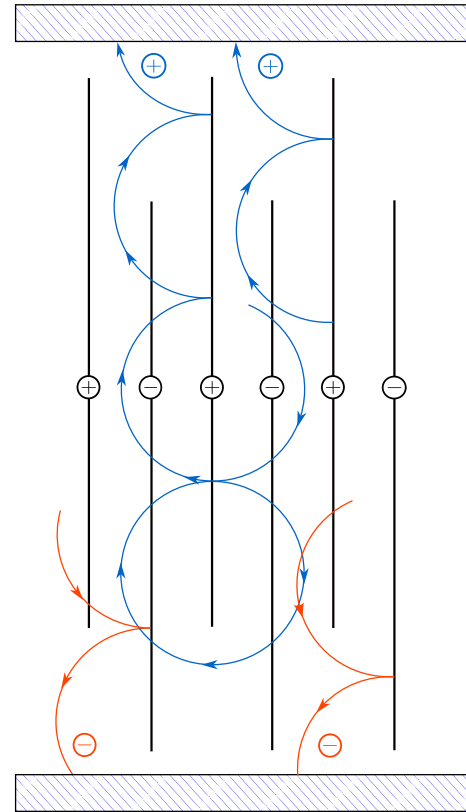


FIG. 6. Sketch of a geometry in order to collect the positive and negative charge carriers at the top and bottom metallic contacts, respectively. The negative potential barriers, which act as reflectors for the electrons, are extended downward, while the positive potential barriers are extended upward.

the arriving holes. The electrons could be collected by an analogous procedure on the other side. A sketch of such a setup is depicted in Fig. 6. Alternatively, one could collect the electrons and holes by additional (negative and positive) potential barriers parallel to the x axis, which terminate the other potential barriers (forming two combs telescoped into each other) and steer the electrons and holes to the contacts.

VI. EXPERIMENTAL PARAMETERS

In order to discuss the order of magnitude of the involved parameters, let us consider an explicit example for an experimental realization. To this end, we envisage exfoliated graphene monolayers that are sandwiched between sheets of hexagonal boron nitride (hBN), which are 30 nm thick. Boron nitride has several advantages. It is a very good insulator with a high dielectric breakdown threshold on the order of 10^9 V/m [40]. Furthermore, it supports high-quality encapsulated graphene, which can exhibit carrier mobilities exceeding $300\,000$ cm²/Vs and mean free paths of several microns, even at room temperature [41].

Since boron nitride is transparent in the optical and infrared regime, it allows us to stack the graphene layers in order to enhance photon absorption (see Fig. 7). With present-day fabrication techniques, it should be possible to

create patches of a lateral size of around $50\text{ }\mu\text{m}$, which can then be placed side by side and contacted in order to cover a larger area.

To further enhance photon absorption, we place a mirror at the top or bottom of the stack (see Fig. 7). On top of the upper boron-nitrate sheet, metal electrodes of, say, 200 nm width and 30 nm height could be placed via lithography [42–44]. They allow us to adjust the desired potential landscape $\phi(x)$ with the spatial resolution of a few hundred nm. Applying gate voltages in the volt regime could generate electric fields exceeding 10^6 V/m, which are superthreshold for magnetic fields of 1 T or below.

Between the electrodes carrying these gate voltages, the chemical potential is supposed to be close to the Dirac point(s) of graphene such that incident photons can create electron-hole pairs. To illustrate one of the advantages of graphene (as a gapless material), we consider charge carriers with an energy of $\omega = 0.1\text{ eV}$, which would be created by the absorption of infrared photons. Applying a magnetic field of $B = 0.1\text{ T}$ (which can be generated by strong permanent magnets), we get a cyclotron radius (5) of around $1\text{ }\mu\text{m}$, which is comparable to the mean free path in graphene [41]. If the mean free path were much smaller than the cyclotron radius, multiple scattering events would deteriorate the desired effect. The distance between two electrodes with positive voltages should not be much larger than the associated diameter of $2\text{ }\mu\text{m}$ (see Fig. 6), in order

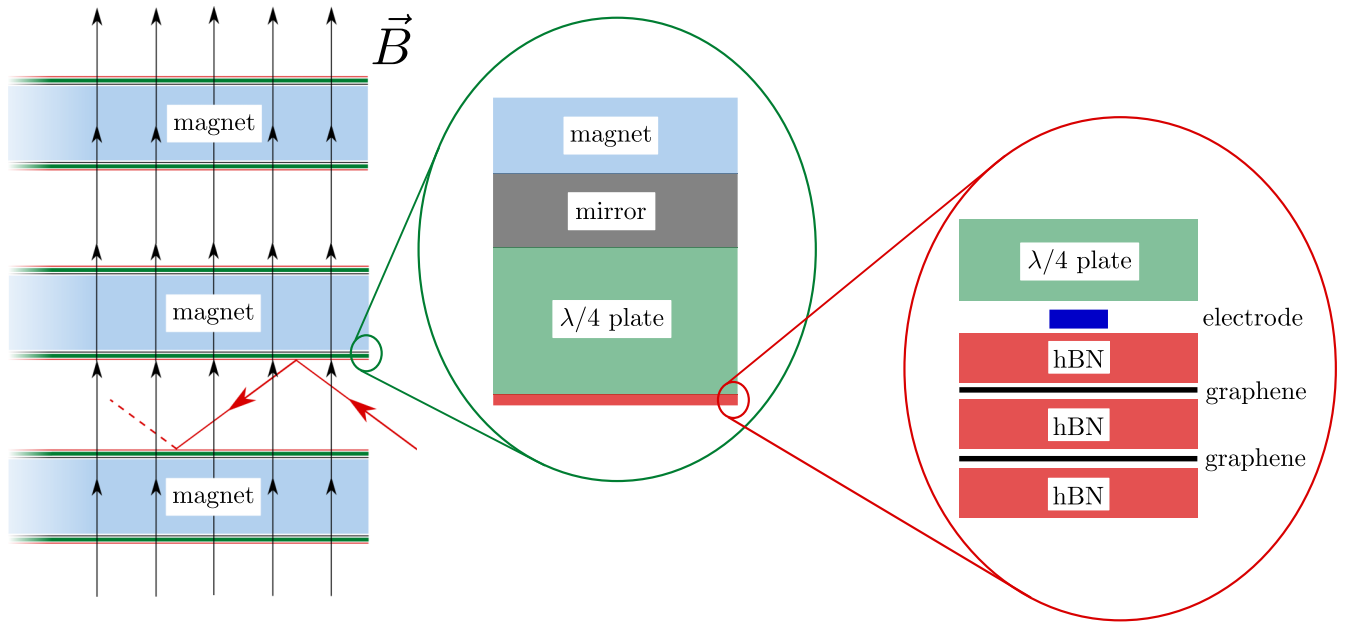


FIG. 7. Sketch (not to scale) of a possible experimental realization. The left figure with a characteristic length scale in the millimeter range shows a possible configuration of permanent magnets (blue) generating the required magnetic field, while the incident optical or infrared radiation (red arrows) is reflected at the mirrors (gray) coating these magnets and thus bounces back and forth between them in order to enable efficient absorption at the graphene layers. The middle figure with a characteristic length scale in the micrometer regime shows a blowup of the coating layer that also contains a transparent $\lambda/4$ plate (green) separating the mirror from the active sheet (red) in order to avoid the suppression of the electric field by the boundary conditions at the mirror. Finally, the right figure depicts the active sheet containing multiple graphene monolayers (black) separated by hexagonal boron nitride (hBN) sheets as well as the electrodes (blue).

to ensure that most of the charge carriers would be guided along the electrodes. Thus, by suitably switching on or off electrodes, one can adapt the geometry to the energy of the charge carriers.

Note that the cyclotron radius (5) is much larger than the Landau length $\ell_B = \sqrt{\hbar/(qB)} \approx 80$ nm as well as the free wavelength $\lambda \approx 90$ nm of the charge carriers, which shows that the semiclassical picture should provide a good approximation. Similarly, the tunneling probability through the potential barriers is exponentially suppressed and can be neglected.

Assuming a high efficiency for photon absorption, which can be achieved by a cavity geometry as in Fig. 7, for example, an incident intensity corresponding to the solar constant of $\mathcal{O}(1 \text{ kW/m}^2)$ could generate a maximum current of $\mathcal{O}(0.1 \text{ A})$ per illuminated area of 1 cm^2 .

The above considerations describe only one example setup. As an alternative, one could, for example, use chemical-vapor-deposition-grown graphene on Si/SiO₂ substrates. Here, the electrode structures can be realized by ion implantation in silicon; the SiO₂ layer acts as the dielectric that insulates the graphene from the gates.

VII. CONCLUSION AND OUTLOOK

Motivated by the potential of graphene for photoelectric or photovoltaic applications, we study the motion of charge carriers in planar graphene under the combined influence of a constant perpendicular magnetic field B and an in-plane electrostatic potential landscape $\phi(x)$, which could be generated, for example, by external electrodes or by doping. To this end, we solve the effective Dirac equation (valid in the vicinity of the Dirac cones) and discuss the associated semiclassical trajectories. We find that by shaping the potential landscape $\phi(x)$ and thereby breaking the $\mathcal{C}$ (charge), $\mathcal{P}$ (parity), and $\mathcal{T}$ (time reversal) symmetries in an appropriate way, we may effectively steer the charge carriers and achieve charge separation. Apart from graphene-based photoelectric or photovoltaic applications, this steering capability could also be exploited for “graphene electronics,” etc.

Since the directed motion of charge carriers in the scenarios considered here already follows from general principles, we expect that the steering phenomenon is quite robust against small perturbations and imperfections, which should help us to realize this effect experimentally. In comparison to scenarios based on geometric constraints (such as graphene folds or edges) discussed previously, see, e.g., Refs. [22,23], the scheme considered here offers several advantages. Since it is based on undisturbed planar graphene, it can be realized and scaled up to larger areas more easily, and it is less prone to defects and imperfections than, for example, a graphene edge. In addition, by adjusting the applied voltages and/or switching on/off

electrodes entirely, one can easily manipulate the potential landscape $\phi(x)$ and thereby adapt it to the intended use (e.g., photovoltaic applications for infrared instead of optical photons).

A. Outlook: Auger effect

One effect that we have not considered here is the Coulomb interaction between the charge carriers, which would require going beyond the single-particle picture in Eq. (1). For example, the Coulomb interaction could cause the Auger-like creation of secondary electron-hole pairs. In this case, a primary charge carrier (generated by photoexcitation, for example) would lose energy (inelastic scattering) by creating a secondary electron-hole pair. For free charge carriers in planar graphene, this process is strongly constrained by energy-momentum conservation. In the presence of imperfections or a graphene edge (or a potential barrier) that can take up the recoil momentum, however, the available phase space is enlarged.

Actually, the measurement of the magnetophotovoltaic effect at graphene edges reported in Ref. [24] indicates such an Auger-like effect, as the primary charge carriers generated by photoexcitation can only account for a fraction (around 1/8) of the measured current. The analytic estimates in Refs. [24] and [22] show that Auger-like effects are expected to be most pronounced at the edge of graphene (or other imperfections such as defects). Since the setup considered here does not involve such sharp edges but smooth potentials $\phi(x)$ instead, Auger-like effects are probably less relevant here, but this should be the subject of further studies.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX: WKB APPROXIMATION

The WKB approach can be employed to derive a semiclassical approximation to the full quantum evolution in Eqs. (1) or (3), which then leads to the dispersion relation (4) and the semiclassical equations of motion (11) and (12). The scheme is very analogous to the WKB

method applied to the stationary Schrödinger equation in one dimension but must be adapted to graphene, i.e., to two-component wave functions Ψ and first-order derivatives ∂_x , see, e.g., Refs. [45–48]. The idea is to split up the full wave function Ψ into a rapidly oscillating phase $e^{iS(x)}$ and a slowly varying amplitude $U(x)$. Thus, the ansatz (2) is generalized via

$$\Psi(t, x, y) = \exp \{-i\omega t +iky + iS(x)\} U(x). \quad (\text{A1})$$

Inserting this WKB ansatz (A1) into Eqs. (1) or (3) and neglecting $\partial_x U$ in comparison to $\partial_x S$, which amounts to replacing ∂_x by $ip_x = i\partial_x S$, we arrive at a linear set of equations for U . To have a nontrivial solution for U , the determinant of the associated matrix must vanish, which yields the dispersion relation (4) relating p_x to k and ω . Thus, the eikonal or phase function $S(x)$ is given by

$$S = \int dx p_x = \int dx \sqrt{\frac{[\omega - q\phi(x)]^2}{v_F^2} - [k + qBx]^2}. \quad (\text{A2})$$

As in the usual Schrödinger equation, the WKB approximation breaks down at the classical turning points where $p_x = 0$. Thus, the full wave Eq. (3) must be used to connect spatially oscillating solutions in the classically allowed region (where p_x is real) across the turning point to the evanescent tail in the classically forbidden region (where p_x is imaginary).

In order to reproduce the semiclassical equations of motion (11) and (12), let us study the time evolution of wave packets. To this end, we consider the superposition of wave functions (A1) for different values of ω and k

$$\Psi(t, x, y) = \int d\omega dk e^{-i\omega t +iky + iS_{\omega k}(x)} U_{\omega k}(x), \quad (\text{A3})$$

where the amplitude function $U_{\omega k}(x)$ is slowly varying in x but peaked in ω and k (centered around the mean values ω_m and k_m). As usual, the peak of the wave packet can be determined by the saddle (or stationary phase) point of the ω and k integrals, which gives the two conditions $\partial S/\partial k = -y$ and $\partial S/\partial \omega = t$. Their differentials $\partial dS/\partial k = -dy$ and $\partial dS/\partial \omega = dt$ with $dS = p_x dx$ yield $p_x dy = p_y dx$ and $p_x v_F dt = |p| dx$, which imply the equation of motion (11).

Note that the mechanical momentum p_y in the y direction is related to the (constant) canonical momentum k via $p_y = k + qBx$, i.e., it does effectively also depend on x . Since k is conserved, the temporal change dp_y/dt is given by $dp_y/dt = qBdx/dt$, which is the y component of the semiclassical Lorentz force (12). The remaining x

component can be obtained from the time derivative of the dispersion relation (4). Using the results above, we find $dp_x/dt = -q\partial_x \phi/v_F - qBdy/dt$, which constitutes the other component of the semiclassical Lorentz force (12).

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