Cyclotron resonance of the magnetic ratchet effect and second harmonic generation in bilayer graphene

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We model the magnetic ratchet effect in bilayer graphene in which a dc electric current is produced by an ac electric field of frequency \( \omega \) in the presence of a steady in-plane magnetic field and inversion-symmetry breaking. In bilayer graphene, the ratchet effect is tunable by an external metallic gate which breaks inversion symmetry. For a tilted magnetic field, the perpendicular field component induces cyclotron motion with frequency \( \omega_c \) and we find that the dc current displays cyclotron resonance at \( \omega_c = \omega \), although this peak in the current is actually smaller than its value at \( \omega_c = 0 \). Second harmonic generation, however, is greatly enhanced by resonances at \( \omega_c = \omega \) and \( \omega_c = 2\omega \) for which the current is generally much larger than at \( \omega_c = 0 \).

I. INTRODUCTION

Recently, there has been interest in the magnetic ratchet effect in two-dimensional electron systems such as semiconductor quantum wells [1,6], graphene [7,9] and bilayer graphene [10]. It is a non-linear effect [11,12] producing a dc electric current in response to ac laser light in the presence of a steady in-plane magnetic field and broken inversion symmetry. Here, we consider the magnetic ratchet in bilayer graphene [10,13,15] for which inversion asymmetry is tunable by applying a gate voltage [14,16–18], and we take into account a tilted magnetic field. If the magnetic field has a perpendicular component, the efficiency of the ratchet should be dramatically increased when the magnetic field has a perpendicular component, the efficiency of the ratchet should be dramatically increased under cyclotron resonance conditions [8] when the cyclotron frequency \( \omega_c \) is close to the ac field frequency \( \omega \). For the dc current, we find that there is indeed a resonance at \( \omega_c = \omega \), but, as a function of \( \omega_c \), the dc current is actually larger at \( \omega_c = 0 \), Fig. 1(a). For second harmonic generation [17–25], however, we find there are resonances at \( \omega_c = \omega \) and \( \omega_c = 2\omega \) which are generally much stronger than the current at \( \omega_c = 0 \), Fig. 1(b).

We begin by summarizing our results. In bilayer graphene, the intraband contribution (which is relevant in the semiclassical regime \( \epsilon_F \gg \hbar \omega \) where \( \epsilon_F \) is the Fermi level) to the dc current density \( \mathbf{J}^{(0)} \) is given by

\[
\mathbf{J}^{(0)} \approx \left\{ (\mathbf{B}_\parallel \times \hat{n}_z) \cdot \mathbf{E}^* + (\mathbf{B}_\parallel \times \hat{n}_z) \cdot \mathbf{E}^* \right\} \mathbf{E} - (\mathbf{B}_\parallel \times \hat{n}_z) |\mathbf{E}|^2 \text{Re}(M_1) + (\mathbf{E} \cdot \mathbf{B}_\parallel \cdot \mathbf{E}^*) + \mathbf{E}^* \cdot (\mathbf{B}_\parallel \cdot \mathbf{E}) - \mathbf{B}_\parallel |\mathbf{E}|^2 \text{Im}(M_1),
\]

for normally-incidence radiation with an in-plane alternating electric field \( \mathbf{E}(t) \) of angular frequency \( \omega \),

\[
\mathbf{E}(t) = \mathbf{E} e^{-i\omega t} + \mathbf{E}^* e^{i\omega t},
\]

and in-plane magnetic field \( \mathbf{B}_\parallel \), \( \hat{n}_z \) is a unit vector in the perpendicular (\( z \)) direction. We include terms that are second order in electric field (i.e. quadratic in electric field amplitudes \( \mathbf{E} \) and \( \mathbf{E}^* \)), and linear in \( \mathbf{B}_\parallel \). Second order (in electric field) effects generally require breaking of spatial inversion symmetry [20,28], and, here, the coefficient \( M_1 \) changes sign upon \( z \rightarrow -z \) inversion.

Eq. (1) describes the contribution arising from a perfectly quadratic dispersion relation \( \epsilon = p^2/2m \) (\( p \) is the magnitude of momentum and \( m \) is mass) when the relaxation rates are independent of energy, and this contribution will generally dominate in bilayer graphene (there will be small corrections when these conditions are not exactly met as described in detail in Section III). Parameter \( M_1 \) describes the response to incoming linearly polarized light and we find that

\[
M_1 \approx -\frac{g e^3 p^2}{16\pi^2\hbar^4} \frac{2(2\Lambda_1 + \epsilon\Lambda_1^*)}{\gamma_{1,0}^2} \left( \frac{1}{\gamma_{1,1}^2} + \frac{1}{\gamma_{-1,1}^2} \right),
\]

where \( g \) is a degeneracy factor (\( g = 4 \) for spin and valley in graphene) and the electronic charge is \( -e \), \( e > 0 \). For a particular material, parameter \( \Lambda_1 \) (which is independent of \( \omega \) and \( \omega_c \)) characterizes the strength of scattering in the presence of an in-plane magnetic field and \( z \rightarrow -z \) asymmetry (we will present the explicit form of \( \Lambda_1 \) for
account for the dependence of the current on electric field frequency \( \omega \) and cyclotron frequency \( \omega_c = eB_1v_y/p \) including the cyclotron resonance effect \( [6] [29] [30] \). Here \( B_1 = |B_1| \) and we consider \( B_\perp = B_1\hat{n}_z \). Also, \( v_y \) is the group velocity \( v_y = de/dp \), so \( \omega_c = eB_1v_y^2/e \) \[31\] for linear dispersion \( \epsilon = v_p \), and \( \omega_c = eB_1/m \) \[32\] for quadratic dispersion \( \epsilon = p^2/(2m) \). Parameters \( \tau_1 \) and \( \tau_2 \) are relaxation times for the first and second angular harmonics of the electronic distribution, respectively (\( \tau_1 \) is the usual momentum relaxation rate that we denote as \( \tau \) from now on).

Figure \[1\,a\] shows the magnitude of the non-linear coefficient \( |M_1| \) as a function of cyclotron frequency \( \omega_c \) for fixed \( \omega \). Although there is a noticeable resonance for \( \omega_c = \omega \) [due to the presence of \( \Gamma_{\pm 1}^{\pm 1} \) in Eq. \[3\]], the ratchet effect is actually strongest for \( \omega_c = 0 \) because of the product \( \Gamma_{0}^{0} \). This product arises because of the need to couple with dc components of the electronic distribution in order to create dc current.

The second harmonic current density \( J^{(2)} \) is given by

\[
J^{(2)} \approx 2Re \left[ 2 \left[ (B_\parallel \times \hat{n}_z) \cdot E \right] E - (B_\parallel \times \hat{n}_z) E^2 \right] N_1 e^{-2i\omega t} \]

\[-2Re \left[ 2 \left[ (B_\parallel \cdot E \right) E - B_\parallel E^2 \right] N_2 e^{-2i\omega t} \right],
\]

where \( E^2 \equiv E \cdot E = E_x^2 + E_y^2 \). This describes the contribution arising from a perfectly quadratic dispersion relation when the relaxation rates are independent of energy (there will be small corrections when these conditions are not exactly met as described in detail in Section \[III\]). These terms describe an effect similar to the Faraday effect \[29\] in that incoherently plane-polarized light results in the emission of a plane-polarized second harmonic with the magnetic field contributing to a rotation of the angle of polarization. Incoherently circularly-polarized light results in circularly-polarized second harmonic generation. The coefficients \( N_1, N_2 \) are given by

\[
N_1 = -\frac{ge^3 p^2}{32\pi^2 \hbar^4} \left( \Lambda_1 + e\Lambda_1' \right) \times \left( \frac{1}{\Gamma_{11}^{11} \gamma_{22} \gamma_{22}^{-1}} + \frac{1}{\Gamma_{11}^{11} \gamma_{22} \gamma_{22}^{-1}} \right),
\]

\[
N_2 = -\frac{ige^3 p^2}{32\pi^2 \hbar^4} \left( \Lambda_1 + e\Lambda_1' \right) \times \left( \frac{1}{\Gamma_{11}^{11} \gamma_{22} \gamma_{22}^{-1}} - \frac{1}{\Gamma_{11}^{11} \gamma_{22} \gamma_{22}^{-1}} \right).
\]

Figure \[1\,b\] shows the magnitude of the non-linear coefficient \( |N_1| \) as a function of cyclotron frequency \( \omega_c \) for fixed \( \omega \) (note that \( |N_2| \) shows almost the same dependence on \( \omega_c \), both qualitatively and quantitatively, except that \( |N_2| = 0 \) for \( \omega_c = 0 \)). In stark contrast to \( |M_1| \), the resonance at \( \omega_c = \omega \) is far stronger than the signal at \( \omega_c = 0 \) because there are no dc components of the electronic distribution involved (i.e. no \( \Gamma^{0}\) factors), and there is also a resonance at \( \omega_c = 2\omega \).

The combinations of electric and magnetic fields in Eqs. \[15\] satisfy spatial symmetries including rotations in the two-dimensional plane \((x,y)\) of the sample. As \( B_\parallel \) is an axial vector, the combination \( B_\parallel \times \hat{n}_z \), appearing in the \( \text{Re}(M_1) \) and \( N_1 \) terms, behaves as a true vector. However, factors in the \( \text{Im}(M_1) \) and \( N_2 \) terms containing \( B_\parallel \) appear to break some reflection symmetries (e.g. reflection in a plane perpendicular to the sample, such as the \( y-z \) plane) but they are in fact satisfied because \( \text{Im}(M_1) \) and \( N_2 \) are odd functions of the combination \( \hat{n}_z \cdot B_\perp \) and, thus, change sign upon such reflections. Hence, \( \text{Im}(M_1) = 0 \) and \( N_2 = 0 \) when \( B_\perp = 0 \) whereas \( \text{Re}(M_1) \) and \( N_1 \) are even functions of \( \hat{n}_z \cdot B_\perp \) and are generally non-zero for \( B_\perp = 0 \).

In the next Section, we describe a phenomenological Drude model in which the ratchet effect may be viewed as arising from the motion of classical particles in the presence of friction created by the in-plane magnetic field. This model correctly accounts for the combinations of fields in Eqs. \[15\], but doesn’t quite account for the frequency dependences in Eqs. \[15\] because it only contains one relaxation rate \( \tau \). However, in Section \[III\] we use the Boltzmann equation to derive general equations for the dc current and second harmonic for a magnetic ratchet in an arbitrary two-dimensional electron system with \( z \rightarrow -z \) asymmetry and an isotropic dispersion \( \epsilon(p) \), including cyclotron motion, too. Section \[IV\] describes scattering in the presence of an in-plane magnetic field in bilayer graphene in order to apply the general equations to it, resulting in the simplified expressions Eqs. \[15\].

II. PHENOMENOLOGICAL DRUDE MODEL

The expressions \[15\] for non-linear current densities may be viewed as being due to the motion of classical particles in the presence of friction created by the in-plane magnetic field. Previously \[1\,3\], a classical, Drude model has been used to provide a simple picture of the origin of ratchet current for intraband transitions. Here, we generalize this approach for the case where the non-linearity is produced by the in-plane magnetic field, and we consider cyclotron motion and second harmonic generation, too. We consider the equation of motion for the average drift velocity \( \mathbf{v}_d \) per electron in a two-dimensional system:

\[
m \frac{d\mathbf{v}_d}{dt} = -e\mathbf{E} - e(\mathbf{v}_d \times \mathbf{B}_\perp) - \frac{m\mathbf{v}_d}{\tau} + \mathbf{F}(\mathbf{v}_d, \mathbf{B}_\parallel),
\]

where \( \tau \) is the momentum relaxation time, and the perpendicular magnetic field \( \mathbf{B}_\perp \) and in-plane ac electric field \( \mathbf{E}(t) \), Eq. \[2\], enter via the Lorentz force. The term \( \mathbf{F}(\mathbf{v}_d, \mathbf{B}_\parallel) \) describes friction due to the presence of the in-plane magnetic field \( \mathbf{B}_\parallel \) which introduces non-linearity into the system. We assume this term is quadratic in
velocity \( \mathbf{v}_d \) and linear in magnetic field \( \mathbf{B}_\parallel \):
\[
\mathbf{F}(\mathbf{v}_d, \mathbf{B}_\parallel) = \alpha \left\{ 2 \left[ (\mathbf{B}_\parallel \times \hat{n}_z) \cdot \mathbf{v}_d - (\mathbf{B}_\parallel \times \hat{n}_z) \right] |\mathbf{v}_d|^2 \right\},
\]
where \( \alpha \) is a phenomenological parameter that characterizes the material properties of the particular system in question. This form of the friction term is obtained by requiring that it behaves as a true vector rather than, say, an axial vector (as \( \mathbf{B}_\parallel \) is an axial vector, the combination \( \mathbf{B}_\parallel \times \hat{n}_z \) behaves as a true vector). In principle, there should be a second phenomenological parameter to characterize the relative weight of the two terms in \( \mathbf{F}(\mathbf{v}_d, \mathbf{B}_\parallel) \), but, for simplicity, we insert its value above [it is determined by demanding that we obtain non-linear contributions with the correct combinations of fields Eqs. (15)].

To solve Eq. (8) we use a harmonic expansion of the velocity,
\[
\mathbf{v}_d(t) = \sum_{\ell = 0, \pm 1, \ldots} \mathbf{v}_d^{(\ell)} e^{-i\omega nt},
\]
which coupled equations for the coefficients \( \mathbf{v}_d^{(\ell)} \):
\[
-i\hbar \omega \mathbf{v}_d^{(\ell)} = -\frac{e}{m} \mathbf{E} \delta_{\ell, 1} - \frac{e}{m} \mathbf{E}^* \delta_{\ell, -1} - \omega_c (\mathbf{v}_d^{(\ell)} \times \hat{n}_z) - \frac{\mathbf{v}_d^{(\ell)}}{\tau} + \frac{\alpha}{m} \sum_n \left\{ 2 \left[ (\mathbf{B}_\parallel \times \hat{n}_z) \cdot \mathbf{v}_d^{(n)} \right] \mathbf{v}_d^{(\ell - n)} - (\mathbf{B}_\parallel \times \hat{n}_z) \left[ \mathbf{v}_d^{(n)} \cdot \mathbf{v}_d^{(\ell - n)} \right] \right\}. \tag{9}
\]

Neglecting the in-plane magnetic field, the linear harmonics are given by
\[
\left( \mathbf{v}_d^{(1)} \right)_x = -\frac{e\tau}{m} \left( \frac{1 - i\omega \tau}{1 - i\omega \tau + (\omega_c \tau)^2} \right) E_x - \omega_c \tau E_y,
\]
\[
\left( \mathbf{v}_d^{(1)} \right)_y = -\frac{e\tau}{m} \left( \frac{1 - i\omega \tau}{1 - i\omega \tau + (\omega_c \tau)^2} \right) E_y + \omega_c \tau E_x,
\]
and \( \mathbf{v}_d^{(-1)} = \left( \mathbf{v}_d^{(1)} \right)^* \). Since current density is related to drift velocity by \( \mathbf{J}(t) = -ne\mathbf{v}_d(t) \) where \( n \) is the electron number density [22], the linear current density may be written as \( \mathbf{J}^{(1)} = -2ne\mathbf{E} \mathbf{v}_d^{(1)} e^{-i\omega nt} \) which yields \( \mathbf{J}^{(1)} = 2\text{Re} \{ n \mathbf{E} \mathbf{v}_d e^{-i\omega nt} \} \) where the conductivity tensor \( \sigma \) has components
\[
\sigma_{xx} = \sigma_{yy} = \frac{(1 - i\omega \tau)\sigma_0}{(1 - i\omega \tau + (\omega_c \tau)^2)}, \tag{10}
\]
\[
\sigma_{xy} = -\sigma_{yx} = -\frac{(1 - i\omega \tau)\sigma_0}{(1 - i\omega \tau + (\omega_c \tau)^2)}, \tag{11}
\]
with dc Drude conductivity \( \sigma_0 = ne^2\tau/m \), as expected [32].

The in-plane magnetic field introduces non-linearity which couples non-linear harmonics to the linear one, Eq. (5). To first order in \( \mathbf{B}_\parallel \), we find that the ratchet current density is given by Eq. (1) and the Drude expression for the non-linear coefficient \( M_1 \) is
\[
M_1^D = -\frac{\alpha e^3}{m^3} \left( \frac{1}{(\Gamma^0, 1)^2} \right) \left( \frac{1}{\Gamma^1, 1} + \frac{1}{\Gamma^{-1, 1}} \right).
\]
Likewise, the second harmonic current density is given by Eq. (2) and the Drude expressions for the non-linear coefficients \( N_1, N_2 \) are given by
\[
N_1^D = -\frac{\alpha e^3}{2m^3} \left[ \frac{1}{(\Gamma^1, 1)^2 \Gamma^2, 1} + \frac{1}{(\Gamma^{-1, 1})^2 \Gamma^{-2, 1}} \right],
\]
\[
N_2^D = \frac{\alpha e^3}{2m^3} \left[ \frac{1}{(\Gamma^1, 1)^2 \Gamma^2, 1} - \frac{1}{(\Gamma^{-1, 1})^2 \Gamma^{-2, 1}} \right].
\]

Comparison with the coefficients derived using the Boltzmann equation Eqs. (13) shows that although the simple Drude model correctly produces the correct combinations of fields Eqs. (11), it doesn’t quite account for the frequency dependence because it only contains one relaxation rate \( \tau \) [the parameter \( \tau_2 \) describing relaxation of the second angular harmonic of the electronic distribution is not included in Eq. (3)]. In order to accurately describe temporal and spatial relaxation of the electronic distribution it is necessary to employ the Boltzmann equation, as described in the next Section.

### III. BOLTZMANN EQUATION

In this section, we derive the intraband contribution to the second order (in electric field) non-linear conductivity due to the presence of an in-plane magnetic field \( \mathbf{B}_\parallel \) for an arbitrary two-dimensional electron system with \( z \to -z \) asymmetry and an isotropic dispersion \( \epsilon(p) \). We consider linear-in-\( \mathbf{B}_\parallel \) terms and we take into account the effect of a perpendicular magnetic field \( \mathbf{B}_\perp \) which introduces cyclotron motion with cyclotron frequency \( \omega_c = eB_\perp v_g/p \). These semi-classical calculations are valid for finite Fermi energy \( \epsilon_F \) with \( \epsilon_F \gg \{ \hbar/\tau, \hbar \omega_c \}, \omega_c \tau \ll 1 \), and we also assume that the electrons are degenerate \( \epsilon_F \gg k_B T \).

We consider a spatially homogeneous system with electron motion described by the Boltzmann equation [32],
\[
-e \left( \mathbf{E}_\parallel + v_g \times \mathbf{B}_\perp \right) \cdot \nabla f(p, t) + \frac{\partial f(p, t)}{\partial t} = S[f], \tag{12}
\]
where the electron distribution function \( f(p, t) \) is a function of momentum \( p \) and time \( t \), \( v_g = v_g (i \cos \phi + j \sin \phi) \) and \( \phi \) is the polar angle of momentum. The in-plane alternating field \( \mathbf{E}_\parallel(t) \), Eq. (2), and the perpendicular magnetic field \( \mathbf{B}_\perp = B_\perp \hat{n}_z \) are accounted for by the Lorentz force, while the in-plane magnetic field \( \mathbf{B}_\parallel \) enters via the the collision integral \( S[f] \),
\[
S[f] = \sum_{p'} \left[ W_{pp'} f(p', t) - W_{p'p} f(p, t) \right]. \tag{13}
\]
Here \( W_{p'p} \) is the scattering rate describing scattering from initial \( p \) to final state \( p' \) in the presence of a scattering potential \( \delta H \), as given by Fermi’s golden rule,
\[
W_{p'p} = \frac{2\pi}{\hbar} |\langle p' | \delta H | p \rangle|^2 \delta(\epsilon_p - \epsilon_{p'}). \tag{14}
\]
We consider static impurities
\[ \delta H = \sum_{j=1}^{N_i} \hat{Y} u_0(x - \mathbf{R}_j), \]
where \( N_i \) is the number of impurities, \( u_0(x - \mathbf{R}_j) \) describes the spatial dependence of the impurity potential. The dimensionless matrix \( Y \) takes account of any additional degree of freedom related to structure within the unit cell, for example, \( A/B \) lattice in graphene. We neglect interference between different impurities and use the Fourier transform of the impurity potential,
\[ \tilde{u}(\mathbf{q}) = \int d^2r \ u_0(x) e^{-i\mathbf{q} \cdot \mathbf{r}/\hbar}. \]

When evaluating the scattering rate \([14]\), we expand the states \( |p, p'\rangle \) in powers of the in-plane magnetic field \( B_y \). For zeroth order, \( B_y = 0 \), we recover the usual relaxation rates for the \( j \)th angular harmonics of the electronic distribution,
\[ \tau^{-1}_{j} = \frac{2\pi}{\hbar} \sum_{p'} |\langle p' | \delta H | p \rangle|^2 \delta(\epsilon_p - \epsilon_{p'}) \times [1 - \cos (j |\phi' - \phi|)], \quad (15) \]
and we write \( \tau_1 = \tau \). To linear order in \( B_y \), we find that the scattering rate may be written generically as
\[ \delta W_{p,p'} = \frac{1}{\hbar^2} |\tilde{u}(p' - p)|^2 \delta(\epsilon_p - \epsilon_{p'}) \times \left\{ \Omega - \Omega_c \cos[2(\phi' - \phi)][B_x (p'_x + p_y) - B_y (p'_x + p_z)] + \Omega_s \sin[2(\phi' - \phi)][B_x (p'_x + p_z) + B_y (p'_y - p_y)] \right\}, \]
where \( \Omega, \Omega_c, \Omega_s \) are angle-independent factors that depend on specific material properties.

In order to solve the Boltzmann equation \([12]\), we use polar coordinates \((p, \phi)\) for momentum and expand the distribution function in terms of \( \phi \) and \( t \) harmonics with coefficients \( f_{m}^{(n)} \):
\[ f(p, t) = \sum_{n,m} f_{m}^{(n)} e^{i m \phi - i n \omega t}, \quad (17) \]
where \( m, n \) are integers. We also perform an harmonic expansion of the impurity potential,
\[ |\tilde{u}(p' - p)|^2 = \sum_{m} \nu_{-m} e^{i m (\phi' - \phi)}, \]
with the constraint that \( \nu_{-m} = \nu_{m} \) as it is an even function of \( (\phi' - \phi) \). Then, we multiply the Boltzmann equation by a factor \( \exp(-i j \phi + i \omega t) \), where \( j, \ell \) are integers, and integrate over a period \( 2\pi \) of angle \( \phi \) and a period of time \( t \). This results in coupling between different harmonic coefficients:
\[ f_{j}^{(\ell)}(\tau_{j}^{-1} - i \ell \omega + i j \omega_c) = \alpha_{j-1} f_{j-1}^{(\ell-1)} + \eta_{j+1} f_{j+1}^{(\ell+1)} \]
\[ + \tilde{\alpha}_{j-1} f_{j-1}^{(\ell-1)} + \tilde{\eta}_{j+1} f_{j+1}^{(\ell+1)} + \delta S_{j}^{(\ell)}, \quad (18) \]
Operators \( \alpha, \eta \) are linear in the electric field,
\[ \alpha_{j} = \frac{e (E_x - i E_y)}{2} \left( \frac{j}{p} + \frac{\partial}{\partial p} \right), \quad (19) \]
\[ \tilde{\alpha}_{j} = \frac{e (E_x - i E_y^*)}{2} \left( -\frac{j}{p} + \frac{\partial}{\partial p} \right), \quad (20) \]
\[ \eta_{j} = \frac{e (E_x + i E_y)}{2} \left( \frac{j}{p} + \frac{\partial}{\partial p} \right), \quad (21) \]
\[ \tilde{\eta}_{j} = \frac{e (E_x + i E_y^*)}{2} \left( -\frac{j}{p} + \frac{\partial}{\partial p} \right), \quad (22) \]
and factors \( \delta S_{j}^{(\ell)} \) account for the linear-in-\( B_y \) correction to scattering, the relevant ones have small values of \( j \):
\[ \delta S_{0}^{(\ell)} = 0, \quad \delta S_{1}^{(\ell)} = \frac{1}{2} \frac{\Gamma(\varepsilon)}{\varepsilon} (B_y - i B_x) \Lambda_{1} f_{1}^{(\ell)}, \quad (23) \]
\[ \delta S_{-1}^{(\ell)} = \frac{1}{2} \frac{\Gamma(\varepsilon)}{\varepsilon} (B_y + i B_x) \Lambda_{1} f_{-1}^{(\ell)}, \quad (24) \]
\[ \delta S_{2}^{(\ell)} = \frac{1}{2} \frac{\Gamma(\varepsilon)}{\varepsilon} [(B_y + i B_x) \Lambda_{1} f_{1}^{(\ell)} + (B_y - i B_x) \Lambda_{2} f_{3}^{(\ell)}], \quad (25) \]
\[ \delta S_{-2}^{(\ell)} = \frac{1}{2} \frac{\Gamma(\varepsilon)}{\varepsilon} [(B_y - i B_x) \Lambda_{1} f_{-1}^{(\ell)} + (B_y + i B_x) \Lambda_{2} f_{-3}^{(\ell)}]. \quad (26) \]

Here, \( \Gamma(\varepsilon) \) is the electronic density of states per spin and per valley, per unit area \( \Gamma = p/(2\pi h^2 v_g) \), parameters \( \Lambda_1, \Lambda_2 \) are given by
\[ \Lambda_1 = \Omega_0 (v_0 - v_2) + \frac{1}{2} \Omega_2 (v_0 - 2v_2 + v_4), \quad (27) \]
\[ \Lambda_2 = \Omega_0 (v_0 + v_1 - v_2 + v_3) + \frac{1}{2} \Omega_2 (v_0 - 2v_2 - v_3 + v_4 + v_5), \quad (28) \]
\[ \Lambda_3 = \Omega_0 (v_0 - v_3 - v_4 + v_5). \quad (29) \]

The current density is given by \([32]\)
\[ \mathbf{J} = \frac{q e}{\hbar^2} \sum_{p} \mathbf{v}_g f(p, t), \]
where \( g \) is a degeneracy factor \((g = 4 \text{ for spin and valley in graphene})\). Owing to the angular factors in \( v_g \), only the first order angular harmonics \((m = \pm 1)\) in the harmonic expansion \([17]\) survive after integrating over all angles \( \phi \). We express the current as a series in temporal harmonics as
\[ \mathbf{J} = \mathbf{J}^{(0)} + \mathbf{J}^{(1)} + \mathbf{J}^{(2)} + \ldots, \quad (30) \]
\[ \mathbf{J}^{(0)} = -\frac{q e}{\hbar^2} \sum_{p} \mathbf{v}_g \left( f_1^{(0)} e^{i \phi} + f_{-1}^{(0)} e^{-i \phi} \right), \quad (31) \]
\[ \mathbf{J}^{(n)} = -\frac{q e}{\hbar^2} \sum_{p} \mathbf{v}_g \left[ \left( f_1^{(n)} e^{i \phi} + f_{-1}^{(n)} e^{-i \phi} \right) e^{-i n \omega t} \right. \]
\[ \left. + \left( f_{-1}^{(n)} e^{i \phi} + f_1^{(-n)} e^{-i \phi} \right) e^{i n \omega t} \right]; \quad n \geq 1. \quad (32) \]

The coupled equations \([13]\) are used to express harmonics \( f_{j}^{(\ell)} \) in terms of the equilibrium distribution \( f_{0}^{(0)} \). Thus, it is possible to calculate each of the harmonic current densities \( \mathbf{J}^{(0)}, \mathbf{J}^{(1)}, \mathbf{J}^{(2)}, \) which we describe below, beginning with the linear response \( \mathbf{J}^{(1)} \).
A. Linear response $J^{(1)}$

The leading contribution to the linear current density $J^{(1)}$ arises from the linear-in-electric field terms in (18) (i.e. $\delta S_j(t)$ is irrelevant) with

$$f^{(1)}_{+ \pm} = \frac{e \tau (E_x \pm i E_y)}{2(1 - i \omega \tau \pm i \omega_c \tau)} \frac{\partial f_0(0)}{\partial p}; \quad f^{(-1)}_{+ \pm} = (f^{(1)}_{+ \mp})^*.$$  

For a degenerate electron gas, $\epsilon_F \gg k_B T$, we find that $J^{(1)} = 2Re(\sigma E e^{-i \omega t})$ where the conductivity tensor $\sigma$ has components as given in Eqs. [10-11] with dc Drude conductivity $\sigma_0 = ge^2 v_p \tau / (4\pi e^2)$ (all parameters are evaluated on the Fermi surface). For a system with quadratic dispersion $\epsilon = p^2 / (2m)$ (such as bilayer graphene), then $\sigma_0 = ge^2 e^2 \tau / (4\pi e^2)$ for linear dispersion $\epsilon = v_p p$ (monolayer graphene), then $\sigma_0 = ge^2 e^2 / (4\pi e^2)$ [31 34 35].

B. Ratchet dc current $J^{(0)}$

The dc current density, Eq. (21), may be written as

$$J^{(0)} = \left\{ \left[ (B_\parallel \times \hat{n}_z) \cdot E \right] E^* + \left[ (B_\parallel \times \hat{n}_z) \cdot E^* \right] \right\} Re(M_1) + \left[ (B_\parallel \times \hat{n}_z) \right] |E|^2 Re(M_2) + iB_\parallel \left[ (E \cdot E^*) \cdot \hat{n}_z \right] Re(M_3) + \left\{ (E \cdot E^*) + E^* (B_\parallel \cdot E) - B_\parallel |E|^2 \right\} Im(M_1) - B_\parallel |E|^2 Im(M_2) + iB_\parallel \left( E \cdot E^* \right) Im(M_3)$$  

(22)

In terms of components, this may be expressed [6] as

$$J_x^{(0)} = B_x (-|E|^2 ImM_2 + \Theta_1 ImM_1 - \Theta_2 ReM_1 + \Theta_3 ReM_3),$$

$$J_y^{(0)} = B_y (-|E|^2 ImM_2 + \Theta_1 ReM_1 + \Theta_2 ImM_1 + \Theta_3 ImM_3),$$

$$J_z^{(0)} = B_z (-|E|^2 ReM_2 + \Theta_1 ReM_1 + \Theta_2 ImM_1 - \Theta_3 ImM_3) + B_y (-|E|^2 ImM_2 - \Theta_1 ImM_1 + \Theta_2 ReM_1 + \Theta_3 ReM_3),$$

where $\Theta_1 = (|E_x|^2 - |E_y|^2), \Theta_2 = (E_x E_y^* + E_y E_x^*), \Theta_3 = i(E_x E_y^* - E_y E_x^*)$ [30].

For a degenerate electron gas, $\epsilon_F \gg k_B T$, we find the three coefficients are given by

$$M_1 = -\frac{ge^3}{32\pi^2 \hbar^4} \left( \frac{1}{\tau_{1,1}} + \frac{1}{\tau_{-1,1}} \right) \times \left[ \frac{4A_1 p^2}{\tau_{0,1} \tau_{0,2}} + v_y p^3 \left( \frac{\Lambda_1}{\tau_{0,1} \tau_{0,2}} \right) \right],$$

$$M_2 = \frac{ge^3 A_1 p^2}{32\pi^2 \hbar^4} \left[ \frac{1}{\tau_{1,1} \tau_{1,1}} + \frac{1}{\tau_{-1,1} \tau_{-1,1}} \right] \times \left[ \frac{1}{\tau_{0,1}} - \frac{p v_y}{\tau_{0,1}} - v_y p \left( \frac{1}{\tau_{0,1}} \right) \right],$$

$$M_3 = i \frac{ge^3 A_1 p^2}{32\pi^2 \hbar^4} \left[ \frac{1}{\tau_{1,1} \tau_{1,1}} + \frac{1}{\tau_{-1,1} \tau_{-1,1}} \right] \times \left[ \frac{1}{\tau_{0,1}} - \frac{p v_y}{\tau_{0,1}} - v_y p \left( \frac{1}{\tau_{0,1}} \right) \right],$$

(23)

where $(\ldots)' \equiv \partial(\ldots) / \partial \epsilon$ and all parameters are evaluated on the Fermi surface. The terms $Re(M_1), Re(M_2), Re(M_3)$ are all even function of $\hat{n}_z \cdot B_\parallel$, whereas $Im(M_1), Im(M_2), Im(M_3)$ are odd functions, thus they are zero for $B_\parallel = 0$. These equations generalize those in Refs. [3 6 8 10] and describe the intraband contribution to the ratchet effect in a two-dimensional material with isotropic dispersion. Parameters such as the scattering times $\tau, \tau_2$, Eq. (15), and $\Lambda_1$, Eq. (19), are specific to the given material, we will describe them for bilayer graphene in Section IV.

The coefficients $M_1, M_2, M_3$ describe the response to different polarizations of light: $M_2$ characterizes the effect of unpolarized light, $M_1$ describes additional terms that appear if the light is linearly polarized, $M_3$ includes additional terms that occur for circular polarization. In particular, for incoming linearly-polarized light, $E_x(t) = E_0 \cos \theta \cos \omega t$ and $E_y(t) = E_0 \sin \theta \cos \omega t$ where $\theta$ is the polarization angle, then

$$E_x^* = E_x = \frac{E_0}{2} \cos \theta; \quad E_y^* = E_y = \frac{E_0}{2} \sin \theta.$$  

(24)

In this case, the current density may be expressed as

$$J^{(0)} = \frac{E_0^2}{4} B_{||}|M_1| \left\{ i \cos(2 \theta - \varphi - \chi_1 + \pi/2) \right. + \left. j \sin(2 \theta - \varphi - \chi_1 + \pi/2) \right\},$$

$$+ \frac{E_0^2}{4} \left( (B_\parallel \times \hat{n}_z) \right. ReM_2 - \left. B_{||} ImM_2 \right),$$

(25)

where $\chi_1 = \arg(M_1)$ and $\varphi$ is the polar angle of the in-plane magnetic field $B_{||} = (B_x, B_y, 0) = B_{||} (cos \varphi, sin \varphi, 0)$ and $B_{||} = ||B_{||}|$. The $M_1$ term produces current in a direction determined by the polarization angle $\theta$, the magnetic field direction $\varphi$ and the phase $\chi_1$ of the $M_1$ coefficient, whereas $M_2$ describes current in a direction solely determined by the parallel field. For unpolarized light, the $M_1$-related current is zero, but the $M_2$ current survives.
For circularly-polarized light $E_x(t) = E_0 \cos \omega t$ and $E_y(t) = \mu E_0 \sin \omega t$, where $\mu = \pm 1$ indicates left- or right-handed polarization, then

$$E_x^* = E_x = \frac{E_0}{2}; \quad E_y^* = -E_y = -i\mu \frac{E_0}{2}.$$  \hspace{1cm} (26)

Then, the dc current density is given by

$$J^{(0)} = \frac{E_0^2}{2} B_\parallel |M_2| \left\{ i \cos (\varphi - \chi_2 - \pi/2) \right. + \left. j \sin (\varphi - \chi_2 - \pi/2) \right\},$$

$$+ \mu \frac{E_0^2}{2} B_\parallel |M_3| \left\{ i \cos (\varphi - \chi_3) + j \sin (\varphi - \chi_3) \right\},$$

where $\chi_2 = \text{arg}(M_2)$, $\chi_3 = \text{arg}(M_3)$, indicating that the direction of the current is determined by the magnetic field direction and the phase of the $M_2$, $M_3$ coefficients.

C. Second-harmonic generation $J^{(2)}$

The second harmonic of the current density, Eq. (21), may be expressed as

$$J^{(2)} = 2 \text{Re} \left\{ \left[ (B_\parallel \times \hat{n}_z) \cdot \mathbf{E} \right] E^2 N_3 + B_\parallel E^2 N_4 \right\} e^{-2i\omega t}$$

$$+ 2 \text{Re} \left\{ \left[ (B_\parallel \times \hat{n}_z) \cdot \mathbf{E} \right] \left[ (B_\parallel \times \hat{n}_z) \cdot \mathbf{E} \right] N_3 e^{-2i\omega t} \right\}$$

$$- 2 \text{Re} \left\{ \left[ (B_\parallel \times \hat{n}_z) \cdot \mathbf{E} \right] \left[ (B_\parallel \times \hat{n}_z) \cdot \mathbf{E} \right] N_3 e^{-2i\omega t} \right\},$$

where $\mathbf{E}^2 \equiv \mathbf{E} \cdot \mathbf{E} = E_x^2 + E_y^2$. In terms of components, this may be written as

$$J^{(2)} = 2 \text{Re} \left\{ \left[ N_1 (B_y \theta_4 - B_x \theta_5) - N_2 (B_x \theta_4 + B_y \theta_5) \right] e^{-2i\omega t} \right\},$$

$$J^{(2)} = 2 \text{Re} \left\{ \left[ N_1 (B_y \theta_4 + B_y \theta_5) + N_2 (B_y \theta_4 - B_x \theta_5) \right] e^{-2i\omega t} \right\},$$

where $\theta_4 = (E^2_x - E^2_y)$, $\theta_5 = 2E_x E_y$ and $\theta_6 = (E^2_x + E^2_y)$. For a degenerate electron gas, $\epsilon_F \gg k_B T$, we find the coefficients $N_i$ are given by

$$N_1 = \frac{ge^3 v_g p}{2\pi \hbar^2} \left[ \frac{1}{\Upsilon_{1,1}} \left( \frac{v_g p \Gamma_1}{\Upsilon_{2,2} \Upsilon_{2,1}} \right) e^{-2i\omega t} + \frac{1}{\Upsilon_{1,-1}} \left( \frac{v_g p \Gamma_1}{\Upsilon_{2,2} \Upsilon_{2,-1}} \right) e^{-2i\omega t} \right],$$

$$N_2 = \frac{-ige^3 v_g p}{2\pi \hbar^2} \left[ \frac{1}{\Upsilon_{1,1}} \left( \frac{v_g p \Gamma_1}{\Upsilon_{2,2} \Upsilon_{2,1}} \right) e^{-2i\omega t} - \frac{1}{\Upsilon_{1,-1}} \left( \frac{v_g p \Gamma_1}{\Upsilon_{2,2} \Upsilon_{2,-1}} \right) e^{-2i\omega t} \right],$$

$$N_3 = \frac{ge^3 v_g p \Gamma_1}{2\pi \hbar^2} \left[ \frac{1}{\Upsilon_{1,1} \Upsilon_{1,2} \Upsilon_{2,1}} + \frac{1}{\Upsilon_{1,-1} \Upsilon_{1,-2} \Upsilon_{2,-1}} \right]$$

$$- \frac{p}{\Upsilon_{1,1} \Upsilon_{1,2}} \left( v_g \right)' e^{-2i\omega t} + \frac{p}{\Upsilon_{1,-1} \Upsilon_{1,-2}} \left( v_g \right)' e^{-2i\omega t} \right],$$

$$N_4 = \frac{ige^3 v_g p \Gamma_1}{2\pi \hbar^2} \left[ \frac{1}{\Upsilon_{1,1} \Upsilon_{1,2} \Upsilon_{2,1}} - \frac{1}{\Upsilon_{1,-1} \Upsilon_{1,-2} \Upsilon_{2,-1}} \right]$$

$$- \frac{p}{\Upsilon_{1,1} \Upsilon_{1,2}} \left( v_g \right)' e^{-2i\omega t} + \frac{p}{\Upsilon_{1,-1} \Upsilon_{1,-2}} \left( v_g \right)' e^{-2i\omega t} \right].$$

Note that coefficients $N_1$, $N_3$ are even functions of $\hat{n}_z \cdot \mathbf{B}$, whereas $N_2$, $N_4$ are odd, thus $N_2 = N_4 = 0$ for $\mathbf{B} = 0$.

The outgoing linearly-polarized light Eq. (24), the current density is

$$J^{(2)} = \frac{E_0^2}{2} \left[ N_1 \cos (2\omega t - \psi_1) \left\{ i \cos (2\theta - \varphi + \pi/2) \right. \right.$$

$$+ \left. j \sin (2\theta - \varphi + \pi/2) \right\},$$

$$+ \frac{E_0^2}{2} B_\parallel N_2 \cos (2\omega t - \psi_2) \left\{ i \cos (2\theta - \varphi + \pi) \right. \right.$$

$$+ \left. j \sin (2\theta - \varphi + \pi) \right\},$$

$$+ \frac{E_0^2}{2} N_3 \left( B_\parallel \times \hat{n}_z \right) \cos (2\omega t - \psi_3),$$

$$+ \frac{E_0^2}{2} N_4 \left| B_\parallel \right| \cos (2\omega t - \psi_4),$$

where the phases $\psi_1 = \text{arg}(N_1)$, $\psi_2 = \text{arg}(N_2)$, $\psi_3 = \text{arg}(N_3)$, $\psi_4 = \text{arg}(N_4)$ describe a time lag between the incoming light and the produced current. For the $N_1$ and $N_2$ terms, the in-plane magnetic field rotates the polarization direction as in the Faraday effect whereas, for the $N_3$ and $N_4$ terms, the outgoing linear polarization direction is solely determined by the parallel field. Note that, for unpolarized light, the $N_1$ and $N_2$ related currents are zero, but the $N_3$ and $N_4$ currents survive.

For incoming circularly-polarized light Eq. (26), the $N_3$ and $N_4$ related currents are zero, and the current density
is
\[
\mathbf{J}^{(2)} = E_0^2 \mathbf{B}_y |N_1| \left\{ i \cos (2 \omega t - \psi_1 - \mu(\varphi - \pi/2)) + \mu \mathbf{J}_x \sin (2 \omega t - \psi_1 - \mu(\varphi - \pi/2)) \right\}, \\
+ E_0^2 \mathbf{B}_y |N_2| \left\{ i \cos (2 \omega t - \psi_2 - \mu \varphi - \pi) + \mu \mathbf{J}_x \sin (2 \omega t - \psi_2 - \mu \varphi - \pi) \right\}.
\]

Thus, the generated current is also circularly polarized, the direction of the magnetic field contributes to the phase lag.

IV. BILAYER GRAPHENE

A. Four-component Hamiltonian

In order to apply the general equations for the second-order conductivities derived in Section III to a particular system, it is necessary to model scattering in the presence of an in-plane magnetic field in that system in order to derive the form of parameters \( \Lambda_1 \) [Eq. (19)] and \( \Omega_c, \Omega_s \) [Eq. (16)]. For bilayer graphene, this has been done previously [10] in order to model the magnetic ratchet, here we also include cyclotron resonance and second harmonic generation. We will briefly describe electronic scattering in bilayer graphene in the presence of an in-plane magnetic field but we refer the reader to [10] for further details.

Bilayer graphene has four atomic sites in the unit cell, Fig. 2(a), we label them as A1, B1 on the lower layer, and A2, B2 on the upper layer. Sites B1 and A2 lie directly below and above each other, and, as a result, their orbitals are relatively-strongly coupled and these sites are referred to as ‘dimer’ sites. We employ a Cartesian coordinate system with the graphene lying in the \( x-y \) plane, \( z \) in the perpendicular direction, the lower layer of the bilayer is at \( z = -d/2 \), the upper layer at \( z = d/2 \), where \( d \) is the interlayer spacing. We use the tight-binding model \([13, 15, 37, 38]\) with one \( p_z \) orbital per site, and we take the in-plane magnetic field into account \( \mathbf{B}_y = (B_x, B_y, 0) \) with a vector potential \( \mathbf{A} = z(B_y, -B_x, 0) \) that preserves translational symmetry in the \( x-y \) graphene plane.

The vector potential enters the model through a line integral appearing in the matrix elements, for example, the matrix element for in-plane hopping between an A atom at \( \mathbf{R}_A \) and three nearest-neighbor \( B \) atoms at \( \mathbf{R}_{B_j}, j = 1, 2, 3 \), is given by

\[
H_{AB} = \gamma_0 \sum_{j=1}^{3} \exp \left( i \mathbf{k} \cdot (\mathbf{R}_{B_j} - \mathbf{R}_A) - \frac{ie}{\hbar} \int_{\mathbf{R}_{B_j}}^{\mathbf{R}_A} \mathbf{A} \cdot d\mathbf{l} \right),
\]

where \( \gamma_0 \) is a tight-binding parameter and \( \mathbf{k} \) is the wave vector. Two non-equivalent valleys are located at the Brillouin zone corners (K points), wave vector \( \mathbf{K}_{\xi} = \xi(4\pi/3a, 0), \xi = \pm 1 \), and, in the vicinity of these points, the in-plane momentum is \( \mathbf{p} = (p_x, p_y, 0) = \hbar \mathbf{K}_{\xi} \). Keeping linear in \( \mathbf{p} \) and linear in \( \mathbf{B}_y \) contributions, the Hamiltonian [10] in a basis of A1, B1, A2, B2 sites is

\[
H = \begin{pmatrix}
U_1 & v_3 \pi_1^\dagger & -v_4 \pi_1^\dagger & -v_3 \pi_1^\dagger \\
v_3 \pi_1 & U_1 + \delta & \gamma_1 & -v_4 \pi_1^\dagger \\
-v_3 \pi_1^\dagger & -\gamma_1 & U_2 + \delta & v_4 \pi_2^\dagger \\
v_4 \pi_2 & v_4 \pi_2^\dagger & -v_4 \pi_2 & U_2
\end{pmatrix},
\]

where \( v = \sqrt{3a\gamma_0/(2\hbar)} \) represents in-plane nearest-neighbour A1-B1, A2-B2 hopping, \( a \) is the lattice constant, \( \gamma_1 \) describes vertical interlayer coupling, \( v_3 = \sqrt{3a\gamma_3/(2\hbar)} \) represents skew interlayer A1-B2 hopping, and \( v_4 = \sqrt{3a\gamma_4/(2\hbar)} \) represents skew interlayer A1-A2, B1-B2 hopping, Fig. 2(a). The on-site energies of the two layers are characterized by \( U_1, U_2 \), while \( \delta \) describes an energy difference between B1, A2 (dimer sites) and A1, B2 (non-dimers) [15, 38, 41]. Complex momentum operators are labeled \( \pi_1 \) for the lower layer, \( \pi_2 \) for the upper layer and \( \pi \) for interlayer hopping:

\[
\pi = \xi p_x + i p_y, \\
\pi_1 = \xi(p_x - b_y) + i(p_y + b_x), \\
\pi_2 = \xi(p_x + b_y) + i(p_y - b_x),
\]

where the magnetic field, written in dimensions of momentum, is \( b_x = e dB_x/2, b_y = e dB_y/2 \).

B. Two-component reduced low-energy Hamiltonian

There are four \( p_z \) orbitals in the unit cell and Hamiltonian [31] describes four energy bands near each K point, two conduction bands, two valence bands, Fig. 2(b). Of these, one of the conduction bands touches one valence
band at the K point with an approximately quadratic dispersion \( \epsilon = v^2p^2/\gamma_1 \) near zero energy \([14\nobreak,15\nobreak,37\nobreak,38]\). The other two bands are split away from the touching point by \( \pm \gamma_1 \) because the orbitals corresponding to the B1, A2 (dimer) sites are strongly coupled by \( \gamma_1 \). Thus, it is possible to represent the electronic behavior at low energy (less than \( |\gamma_1| \)) where there are only two bands by eliminating the components in Hamiltonian \((31)\) related to B1, A2 (dimer) sites, resulting in a two-component Hamiltonian for the A1, B2 (non-dimer) sites. This process has been described before \([13\nobreak,15]\), and including the in-plane magnetic field \([10]\), the two-component Hamiltonian in an A1, B2 basis is

\[
H = -\frac{v^2}{\gamma_1} \begin{pmatrix} 0 & (\pi^1)^2 \\ \pi^2 & 0 \end{pmatrix} + \frac{\Delta}{2} \begin{pmatrix} 1 - 2v^2p^2/\gamma_1 & 1 & 0 \\ 0 & -1 \end{pmatrix} - \frac{2v^2}{\gamma_1} \begin{pmatrix} v_4 + \delta \gamma_1 & 0 \\ \gamma_1 & 0 \end{pmatrix} \begin{pmatrix} p \times b \end{pmatrix}_2 \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} - \frac{v^4\Delta}{\gamma_1} \begin{pmatrix} 0 & i\pi^1\beta^1 \\ -i\pi^1\beta^1 & 0 \end{pmatrix},
\]

(32)

where \( \beta = b_x + i\xi b_y \), \( \beta^1 = b_x - i\xi b_y \), \( p = |p| \), and interlayer asymmetry is \( \Delta = U_1 - U_2 \). To derive Eq. (32), we neglected a number of contributions, further details may be found in Ref. \([10]\). We excluded terms that are quadratic or higher in the magnetic field, cubic or higher in \( v^2p^2/\gamma_1 \) and cubic or higher in other small parameters \( v_4/\gamma_1 \), \( \delta/\gamma_1 \), \( U_1/\gamma_1 \) and \( U_2/\gamma_1 \). Note that tight-binding models that neglect parameters \( v_4 \) and \( \delta \) omit the linear-in-\( B_|| \) terms \([12\nobreak,38]\). We also excluded \( \gamma_3 \) because it doesn’t lead to magnetic field dependent terms, thus, it will only create small cross terms in the scattering probability, which we neglect. Additionally, we also neglect terms that are proportional to the unit matrix in \((A1, B2)\) space because they don’t influence electronic scattering despite having a small effect on the dispersion relation.

The first term in the Hamiltonian \((32)\) represents chiral quasiparticles \([13\nobreak,14]\) with the pseudospin direction \([14]\) the relative amplitude of the electronic wave function in the lattice \((A1, B2)\) space in the graphene plane and fixed to the direction of the electronic momentum \( p = (p_x, p_y, 0) \). This term results in a quadratic dispersion \( \epsilon = v^2p^2/\gamma_1 \), and we assume that the other terms in \((32)\) are a small perturbation with respect to this dominant one. The second term in \((32)\) describes a gap in the spectrum \([14]\) due to different energies \( U_1, U_2 \) on the two layers characterized by interlayer asymmetry \( \Delta = U_1 - U_2 \). Such interlayer asymmetry could be induced by an external gate voltage and, thus, parameter \( \Delta \) is, in principle, tunable. Also, the presence of a substrate could break inversion symmetry by creating a different electrostatic potential in one layer of the bilayer as compared to the other. Substrate-induced \( \Delta \) could be significant in certain circumstances: for example, it has been estimated to be \( 30 \text{ meV} \) for rippled graphene on SiO\(_2\) \([19]\) and for graphene on hexagonal boron nitride at a small misalignment angle \([50\nobreak,52]\).

The in-plane magnetic field appears in two different terms. The first \([\text{third term in } (32)]\) arises due to small, intrinsic lattice parameters \( v_4 \) and \( \delta \), and it tends to open a gap with a direction dependent on the \( z \) component of the Lorentz force \( p \times B_\parallel \). The magnetic field also appears in the fourth term in \((32)\), this term only appears when there is non-zero interlayer asymmetry \( \Delta \). Next, we use Hamiltonian \((32)\) to calculate the correction to the scattering rate \([14]\) due to the in-plane field.

C. Electron scattering

We determine the scattering rate \([14]\) using Hamiltonian \((32)\) to determine eigenstates \( |p\rangle \) and \( |p'\rangle \) in the presence of the in-plane field \( B_\parallel \), with scattering \( \delta H \) caused by static impurities \((15)\). The dimensionless matrix \( \hat{Y} \) in \( \delta H \) accounts for structure in the A1, B2 lattice degrees of freedom, that is, a possible asymmetry in the level of disorder on the two layers of the bilayer. As representative examples, we consider disorder that is symmetric with equal amounts of scattering on the two layers, \( \hat{Y} = I \) where \( I \) is the unit matrix, and we consider asymmetric disorder, \( \hat{Y} = (I + \delta \hat{Y})/2 \), with scattering limited to the lower \((\zeta = 1)\) or upper \((\zeta = -1)\) layer. Thus, the linear-in-\( B_\parallel \) part of the scattering rate may be written \([10]\) in a general form as in Eq. \((16)\) where the angle-independent factors \( \Omega, \Omega_s, \Omega_a \) are:

\[
\Omega^{(s)}(\epsilon) = \frac{\pi edn\Delta \gamma_1}{2\hbar v^2p^4} \left( \frac{\gamma_4}{\gamma_0} + \frac{\delta}{\gamma_1} \right) \left( 1 - \frac{2v^2p^2}{\gamma_1^2} \right),
\]

\[
\Omega^{(s)}(\epsilon) = \frac{\pi edn\Delta \gamma_4}{2\hbar p^2\gamma_1^2},
\]

(33)

for symmetric disorder, and for asymmetric disorder:

\[
\Omega^{(a)}(\epsilon) = \frac{\pi edn\gamma_4}{2\hbar p} \left( \frac{\gamma_4}{\gamma_0} + \frac{\delta}{\gamma_1} \right) \left( s\zeta - \Delta + \frac{\Delta \gamma_1}{2v^2p^2} \right),
\]

\[
\Omega^{(a)}(\epsilon) = \Omega^{(a)} = 0,
\]

(34)

where the density of impurities is \( n_i = N_i/L^2 \), and \( s = +1 \) \((s = -1)\) for states in the conduction (valence) band. Clearly, for symmetric disorder, there must be interlayer asymmetry \( \Delta = U_1 - U_2 \) to create \( B_\parallel \)-dependent terms, but, for asymmetric disorder, this is not necessary.

D. Ratchet dc current \( J^{(0)} \) in bilayer graphene

For bilayer graphene, we assume that the dispersion is quadratic \([14]\) \( \epsilon = v^2p^2/\gamma_1 \equiv \nu^2/(2m) \) with mass \( m = \gamma_1/2v^2 \) and \( v_3 = 2v^2/\gamma_1 \). Then, the factor \( 1 - p_\nu^2 - \nu_0 p^2 \) in coefficients \( M_2 \) and \( M_3 \), Eq. \((23)\), simplifies as \(-\nu_0 p^2 \) and these coefficients are zero unless \( \tau \) is energy dependent. Overscreened Coulomb impurities in bilayer graphene act like short-range scatterers \([33]\), \( u(r - R_j) = \)
\[ u_0 \delta(\mathbf{r} - \mathbf{R}_j) \text{ and } \tilde{a}(\mathbf{p}' - \mathbf{p}) = u_0, \text{ and the scattering rates Eq. (15) are} \]

- symmetric disorder: \[ \tau^{-1} = 2\tau_2^{-1} = \frac{n_1 u_0^2 \gamma_1}{4 \hbar^2 v_F^2}, \]
- asymmetric disorder: \[ \tau^{-1} = \tau_2^{-1} = \frac{n_1 u_0^2 \gamma_1}{8 \hbar^2 v_F^2}. \]

Then, if \( u_0 \) is independent of energy, so is \( \tau \) and \( M_2 = M_3 = 0 \). The potential is isotropic so the only non-zero harmonic is \( u_0 = u_0^2 \) and parameter \( \Lambda_1 \), Eq. (19), becomes \( \Lambda_1 = u_0^2 \Omega + (\Omega_c + \Omega_s)/2 \). Then, the non-linear coefficients, \( M_1^{(a)} \) for asymmetric disorder Eq. (34), are

\[ M_1^{(a)} = \frac{e^4 d}{8\pi \hbar^2 m^2 \gamma_1} \left( \frac{5\gamma_1}{\gamma_0} + \frac{6\delta}{\gamma_1} \right) \times \frac{1}{\Upsilon_{0.1} \Upsilon_{0.2}} \left( \frac{1}{\Upsilon_{1.1}} + \frac{1}{\Upsilon_{1.1}} \right), \]

\[ M_1^{(s)} = -\frac{e^4 d}{2\pi \hbar^2 m^2 \gamma_1} \left( \frac{\gamma_1}{\gamma_0} + \frac{\delta}{\gamma_1} \right) \times \frac{1}{\Upsilon_{0.1} \Upsilon_{0.2}} \left( \frac{1}{\Upsilon_{1.1}} + \frac{1}{\Upsilon_{1.1}} \right), \]

which generalize the results for \( \omega_c = 0 \) from Ref. [10]. Note that both coefficients change sign under \( z \rightarrow -z \) inversion because of the presence of \( \Delta \) (interlayer asymmetry) and \( \zeta \) (asymmetric disorder), and either mechanism (interlayer asymmetry or asymmetric disorder) produces non-zero \( M_1 \).

**E. Second-harmonic generation \( J^{(2)} \) in bilayer graphene**

For quadratic dispersion and energy-independent scattering rates, \( N_3 = N_4 = 0 \) and the second harmonic is given by Eq. (7) with

\[ N_1^{(s,a)} = \frac{e^4 d \Delta}{8\pi \hbar^2 m^2 \gamma_1} \left( \frac{\gamma_4}{\gamma_0} + \frac{\delta}{\gamma_1} \right) \times \left( \frac{1}{\Upsilon_{1.1} \Upsilon_{2.2} \Upsilon_{2.1}} + \frac{1}{\Upsilon_{1.1} \Upsilon_{2.2} \Upsilon_{2.1}} \right), \]

\[ N_2^{(s,a)} = \frac{ie^4 d \Delta}{8\pi \hbar^2 m^2 \gamma_1} \left( \frac{\gamma_4}{\gamma_0} + \frac{\delta}{\gamma_1} \right) \times \left( \frac{1}{\Upsilon_{1.1} \Upsilon_{2.2} \Upsilon_{2.1}} - \frac{1}{\Upsilon_{1.1} \Upsilon_{2.2} \Upsilon_{2.1}} \right), \]

where the numerical factor is \( c^{(s)} = 3/16 \) for symmetric disorder, Eq. (33), and \( c^{(a)} = 1/8 \) for asymmetric disorder, Eq. (34). For the case of asymmetric disorder, there is no contribution that is independent of \( \Delta \), i.e. asymmetric disorder on its own cannot produce the second harmonic in bilayer graphene, in contrast to the dc current \( M_1^{(a)} \).

As discussed in Section [III C], the \( N_1 \) and \( N_2 \)-related currents are zero for incoming unpolarized light, whereas, for incoming linear polarization, Eq. (29), the in-plane magnetic field rotates the polarization direction as in the Faraday effect [25]. For incoming circularly-polarized light, Eq. (30), the generated current is also circularly polarized and the direction of the magnetic field contributes to the phase lag.

**F. Discussion**

The magnitude of the ratchet current in bilayer graphene may be estimated using values of tight-binding parameters determined by infrared spectroscopy [10] (see also Ref. [53]) such as \( \gamma_0 = 3.0 \text{ eV} \), \( \gamma_1 = 0.4 \text{ eV} \), \( \gamma_4 = 0.015 \text{ eV} \), \( \delta = 0.018 \text{ eV} \). We also use interlayer spacing \( d \approx 3.3 \text{ Å} \) and mass \( m \approx 0.05m_e \), where \( m_e \) is the free electron mass. For a typical value \( \tau = 0.15 \text{ ps} \) [54] and with values \( |E| = 10 \text{ keV cm}^{-1} \), \( |B| = 7 \text{ T} \), \( \omega = 2.1 \times 10^{12} \text{ rad s}^{-1} \) from recent experiment [8], we estimate that the ratchet current density for asymmetric disorder Eq. (30) is of the order of \( |J| \sim |M^{(a)}| |B| |E|^2 \sim 1 \text{ mA cm}^{-1} \). The second harmonic, Eqs. (37, 38), is generally of the same order of magnitude, but with an additional small parameter \( \Delta/\epsilon_F \), its precise value is not fixed because \( \epsilon_F \) and \( \Delta \) are both tunable.

Our results apply to intraband transitions in the semiclassical regime \( \epsilon_F \gg \omega \) (note that recent papers [17, 18] consider second harmonic generation in bilayer graphene at higher frequencies). For Fermi energy \( \epsilon_F = 100 \text{ meV} \), say, the frequency corresponding to \( \hbar \omega = 100 \text{ meV} \) is \( \omega \approx 150 \times 10^{12} \text{ rad s}^{-1} \) or linear frequency \( f \approx 24 \text{ THz} \). For \( \omega_c = eB_{\perp}/m \) with \( m \approx 0.05m_e \) for bilayer graphene, where \( m_e \) is the free electron mass, then a perpendicular field \( B_{\perp} = 1 \text{ T} \) corresponds to \( \omega_c \approx 3.5 \times 10^{12} \text{ rad s}^{-1} \). Then, the cyclotron resonance condition \( \omega = \omega_c \) corresponds to linear frequency of light \( f \approx 0.56 \text{ THz} \) (the corresponding energy scale is \( h\omega \approx 15 \text{ meV} \)) which is well within the semiclassical regime considered here.

**V. CONCLUSIONS**

As detailed in Section [III C] we have determined the dc current, Eq. (22), and the second harmonic generation, Eq. (28), for the magnetic ratchet in the semiclassical regime \( \epsilon_F \gg h\omega \) in a two-dimensional electron system. These results apply to systems with arbitrary, isotropic dispersion and energy-dependent scattering rates. For the particular case of bilayer graphene, we assume a perfectly quadratic dispersion relation \( \epsilon = p^2/2m \) and relaxation rates that are independent of energy to produce simplified expressions for the dc current, Eqs. (1, 35, 36), and second harmonic generation, Eqs. (35, 37, 38). We take into account inversion symmetry breaking by disorder and by interlayer asymmetry, the latter may potentially by induced using an external gate and is thus tunable.
In the presence of a tilted field, we find that the dc current has a resonance at \( \omega_c = \omega \) but that the current value is actually largest at \( \omega_c = 0 \). For the second harmonic, however, resonances at \( \omega_c = \omega \) and \( \omega_c = 2\omega \) generally produce currents significantly greater than that at \( \omega_c = 0 \).

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[36] Coefficients \( M_1, M_2, M_3 \) in Eq. (22) are related to \( M_0, M_L, M_C \) in [1] by \( M_1 = M_L, M_3 = M_0, M_3 = iM_C \). These choices are made to be consistent with earlier works [3, 10] and to ensure the coefficients are real for \( \omega_c = 0 \).