Corrigendum: Laser-assisted spin-polarized transport in graphene tunnel junctions

2012 J. Phys.: Condens. Matter 24 266003

K H Ding¹,², Z G Zhu¹ and J Berakdar¹

¹ Institut für Physik, Martin-Luther-Universität Halle-Wittenberg, D-06099 Halle (Saale), Germany
² Department of Physics and Electronic Science, Changsha University of Science and Technology, Changsha, 410076, People’s Republic of China

E-mail: zhengang.zhu@physik.uni-halle.de

Received 11 July 2012
Published 31 July 2012
Online at stacks.iop.org/JPhysCM/24/349501

In the figures the unit of the electric field strength $E_0$ should read kV cm$^{-1}$. We thank M W Wu for drawing our attention to this typographical error.
Laser-assisted spin-polarized transport in graphene tunnel junctions

Kai-He Ding\textsuperscript{1,2}, Zhen-Gang Zhu\textsuperscript{1} and Jamal Berakdar\textsuperscript{1}

\textsuperscript{1} Institut f"{u}r Physik, Martin-Luther-Universit"{a}t Halle-Wittenberg, 06099 Halle (Saale), Germany
\textsuperscript{2} Department of Physics and Electronic Science, Changsha University of Science and Technology, Changsha, 410076, People\textquotesingle{}s Republic of China

E-mail: zhengang.zhu@physik.uni-halle.de

Received 19 March 2012, in final form 28 April 2012
Published 7 June 2012
Online at stacks.iop.org/JPhysCM/24/266003

Abstract

The Keldysh nonequilibrium Green\textquotesingle{}s function method is utilized to theoretically study spin-polarized transport through a graphene spin valve irradiated by a monochromatic laser field. It is found that the bias dependence of the differential conductance exhibits successive peaks corresponding to the resonant tunneling through the photon-assisted sidebands. The multi-photon processes originate from the combined effects of the radiation field and the graphene tunneling properties, and are shown to be substantially suppressed in a graphene spin valve which results in a decrease of the differential conductance for a high bias voltage. We also discuss the appearance of a dynamical gap around zero bias due to the radiation field. The gap width can be tuned by changing the radiation electric field strength and the frequency. This leads to a shift of the resonant peaks in the differential conductance. We also demonstrate numerically the dependences of the radiation and spin valve effects on the parameters of the external fields and those of the electrodes. We find that the combined effects of the radiation field, the graphene and the spin valve properties bring about an oscillatory behavior in the tunnel magnetoresistance, and this oscillatory amplitude can be changed by scanning the radiation field strength and/or the frequency.

(Some figures may appear in colour only in the online journal)

1. Introduction

Laser fields have been demonstrated to be a powerful tool for the exploration and modification of material properties. For example, in conventional semiconducting nanostructures a variety of applications have been established such as various forms of photo-electronic devices (e.g., radiation-controlled field-effect transistors, photodiodes, and light-emitting diodes \cite{1--5}). On the other hand, graphene with its many unusual physical properties (to mention but a few, the half integer quantum Hall effect \cite{6}, Klein tunneling \cite{7} and its conductance properties \cite{8}) offers a new platform to explore radiation effects. The influence of electromagnetic fields on graphene has recently been a subject of intense research. Particularly interesting examples are the linear response to and the frequency dependence of the conductivity \cite{9--11}, the photon-assisted transport \cite{12, 13}, the microwave and far-infrared response \cite{14--16}, the plasmon spectrum \cite{17--19}, and the nonlinear response to electromagnetic radiation \cite{20--23}. The potential applications of graphene in terahertz electronics were pointed out in \cite{24}.

The resonant interaction between graphene and an electromagnetic field may open a dynamical gap in the quasiparticle spectrum of graphene \cite{25}, which leads to a strong suppression of the quasiparticle transmission through a graphene p–n junction. However, a directed current without applying any dc bias voltage can be generated in certain conditions as a result of inelastic quasiparticle tunneling assisted by one- or two-photon absorptions \cite{26}. Another way to generate a direct current is to use shaped, time-delay tuned electromagnetic pulses \cite{23} which also may generate a valley current. Therefore, it seems possible to control the transport properties of diverse graphene tunneling structures by variation of the strength and frequency of the external radiation fields. On the other hand, the potential of graphene...
for spin-dependent transport (spintronic) applications is well
documented by now [27–35]. Motivated by these facts, in
this work we study theoretically the spin-polarized transport
through a graphene spin valve device in the presence of a
monochromatic laser field. The method is based on the
standard Keldysh nonequilibrium Green’s function approach,
as described in [36, 37]. We find that additional peaks emerge
in the bias dependence of the differential conductance that
reflect the resonant tunneling through the photon-induced
sidebands. The strong suppression of the multiple photon
processes stemming from the combined effect of the radiation
field and the tunneling through the graphene leads to a
strong decay of the differential conductance at a high bias
voltage. The resonant interaction of the quasiparticle in
graphene with the radiation field results in a dynamical
gap in the quasiparticle spectrum. When the bias voltage lies
inside this gap, the differential conductance displays a
zero value region situated symmetrically around zero bias.
The width of this region can be tuned by changing the
radiation strength and the frequency, which also cause a
shift of the resonant peaks in the differential conductance.
The tunnel magnetoresistance (TMR) versus the bias voltage
exhibits an oscillatory behavior, and its oscillatory amplitude
can be controlled by the radiation field strength and/or the
frequency.

2. Theoretical model

We consider a spin valve device consisting of an extended
graphene sheet (that defines the x–y plane) contacted by two
ferromagnetic electrodes. A gate voltage $V_g$ applied on the
graphene shifts the Dirac point away from zero energy. The
two electrodes are voltage-biased with respect to each other
with a bias $V$. The electrical current flows in the $x$ direction.
Additionally, we assume that a laser field is homogeneously
irradiating the structure. Within the metallic electrodes we
assume the field to be shielded and thus ignore its effect on
the electrodes. The laser field is monochromatic and is linearly
polarized along the y direction. The low-energy Hamiltonian
for the graphene around the Dirac points has the minimal
coupling form

$$H_G = v_F \sigma \cdot \mathbf{k} - e v_F A(t) \sigma_y, \quad (1)$$

where $\mathbf{k}$ is the operator of the electron momentum in the
graphene plane, $e$ is the electron charge, $\sigma$ is the vector built
out of the Pauli matrices in the sublattice space, and $A(t)$ is
the laser’s vector potential. In the tight-binding description [38],
$v_F = \frac{E_0}{\omega_0}$ with $t_0$ being the nearest-neighbour hopping energy,
and $a$ being the carbon–carbon distance. The vector potential
$A(t)$ is taken as

$$A(t) = \frac{E_0}{\omega_0} \sin \omega_0 t, \quad (2)$$

where $\omega_0$ is the frequency of the radiation field, and $E_0$ is its
amplitude. Diagonalizing the Hamiltonian (1) in the absence
of the radiation field, one finds the eigenvalues

$$\epsilon_{\pm k} = s v_F |\mathbf{k}| + V_g, \quad (3)$$

with $s = \pm$ denoting the band index, and the eigenstates

$$\psi_{\pm k}(\mathbf{r}) = \frac{1}{\sqrt{2}} e^{i \mathbf{k} \cdot \mathbf{r}} a_{\pm k}^\dagger, \quad (4)$$

where $\Omega$ is the volume of the system, and

$$a_{\pm k}^\dagger = \frac{\sqrt{2}}{2} \begin{pmatrix} 1 \\ i \chi_0 \phi(k) \end{pmatrix}, \quad \tan \phi(k) = \frac{k_y}{k_x}. \quad (5)$$

We introduce the field operators

$$\psi_{\tau}(\mathbf{r}, t) = \sum_\mathbf{k} a_{\pm \mathbf{k} \tau} \psi_{\pm \mathbf{k}}(\mathbf{r}), \quad (6)$$

where $a_{\pm \mathbf{k} \tau}$ is the usual annihilation operator for an electron
in the band $s$, with the momentum $\mathbf{k}$ and the spin $\tau$, and then
express the Hamiltonian (1) in the second quantized form as

$$H_G = \sum_{\mathbf{k} \tau} \epsilon_{\pm \mathbf{k}} a_{\pm \mathbf{k} \tau}^\dagger a_{\pm \mathbf{k} \tau} - e v_F A(t) \sum_{\mathbf{k} \tau} d_{\sigma' \mathbf{k}}^\dagger a_{\pm \mathbf{k} \tau} a_{\pm \mathbf{k} \tau}, \quad (7)$$

where

$$d_{\sigma' \mathbf{k}} = \frac{i}{2} \begin{pmatrix} \chi_0 e^{-i \phi(k)} - s' e^{i \phi(k)} \end{pmatrix}. \quad (8)$$

Applying the rotating wave approximation [39], we neglect
the energy nonconserving terms. Accounting for the coupling
between the graphene and the two ferromagnetic electrodes
we write for the Hamiltonian of the complete tunnel junction
in the presence of the laser

$$H = H_G + H_L + H_R + H_T, \quad (9)$$

where

$$H_G = \sum_{\mathbf{k} \tau} \epsilon_{\pm \mathbf{k}} a_{\pm \mathbf{k} \tau}^\dagger a_{\pm \mathbf{k} \tau} - e v_F \sum_{\mathbf{k} \tau} \left\{ A_0 e^{i \omega_0 t} d_{\pm \mathbf{k} \tau} + a_{\pm \mathbf{k} \tau}^\dagger a_{\pm \mathbf{k} \tau} \right\}, \quad (10)$$

and

$$A_0 = -i \frac{\omega_0}{2 \omega_0}, \quad$$

with $0 < \omega_0 \ll \omega_0$. Therefore

$$H_L = \sum_{\mathbf{q} \lambda} \epsilon_{\mathbf{q} \lambda} c_{\mathbf{q} \lambda}^\dagger c_{\mathbf{q} \lambda}, \quad \lambda = L, R, \quad (11)$$

$$H_T = \sum_{\mathbf{q} \lambda \tau} T_{\mathbf{q} \lambda} c_{\mathbf{q} \lambda}^\dagger a_{\pm \mathbf{k} \tau} a_{\pm \mathbf{k} \tau} + H.c. \quad (12)$$

Equations (11) and (12) describe respectively the $\lambda$ electrode
and the coupling between the graphene and the electrodes.
$\epsilon_{\mathbf{q} \lambda}$ is the single electron energy, $c_{\mathbf{q} \lambda}^\dagger (c_{\mathbf{q} \lambda})$ is the
usual creation (annihilation) operator for an electron with
momentum $\mathbf{q}$ and spin $\tau$ in the $\lambda$ electrode.

By introducing a unitary transformation

$$U = \exp \left[ -i \frac{\omega_0 t}{2} \sum_{\mathbf{k} \tau} \left\{ a_{\pm \mathbf{k} \tau}^\dagger a_{\pm \mathbf{k} \tau} - a_{\pm \mathbf{k} \tau} a_{\pm \mathbf{k} \tau}^\dagger \right\} \right], \quad (13)$$
we redefine the Hamiltonian of the system in the rotating reference frame as
\[
\hat{H} = U^{-1}HU + \frac{i}{\hbar} \frac{dU^{-1}}{dt} U
\]
\[
= \sum_{\mathbf{k} \tau} \epsilon_{\mathbf{k}} a_{\mathbf{k} \tau}^\dagger a_{\mathbf{k} \tau} + \sum_{\mathbf{k} \tau} \Delta (a_{\mathbf{k} \tau}^\dagger a_{-\mathbf{k} \tau} + a_{-\mathbf{k} \tau}^\dagger a_{\mathbf{k} \tau})
\]
\[
+ \sum_{\mathbf{q} \mathbf{k} \tau} \epsilon_{\mathbf{q} \tau} c_{\mathbf{q} \tau}^\dagger c_{\mathbf{q} \tau} + \sum_{\mathbf{q} \mathbf{k} \tau} (T^{\mathbf{q} \mathbf{k}}(t) c_{\mathbf{q} \tau}^\dagger a_{\mathbf{k} \tau} + \text{H.c.}),
\]
where
\[
\Delta = \frac{eV Y_0}{2\omega_0}, \quad T^{\mathbf{q} \mathbf{k}}(t) = T_{\mathbf{q} \mathbf{k}, \tau} e^{-i\omega_0 t},
\]
and
\[
\tilde{\epsilon}_{\mathbf{k}} = \epsilon_{\mathbf{k}} - \omega_0/2.
\]
In the calculation of equation (14), we have assumed (as in [25, 26, 40]) that the most important contributions stem from an almost one-dimensional electron motion ($k_x \gg k_y$) in the interaction with the radiation field.

The electric current of the system can be calculated by the time evolution of the occupation number operator of the left electrode,
\[
I = e(N_{L}^-) = \frac{ie}{\hbar} \langle [\hat{H}, N_{L}^-] \rangle,
\]
where
\[
N_{L}^- = \sum_{\mathbf{q} \tau} c_{\mathbf{q} \tau}^\dagger c_{\mathbf{q} \tau}.
\]
Using the nonequilibrium Green’s function method, equation (15) can be further expressed as
\[
I = -\frac{ie}{\hbar} \sum_{\mathbf{q} \tau} \int dt \int \frac{d\epsilon}{2\pi} \text{Tr} \{[G^<_{\mathbf{q} \tau}(t, t_1)] f_\epsilon(\epsilon) G^>_{\mathbf{q} \tau}(t, t_1)\} \Gamma_L,
\]
\[
\times e^{-i(\epsilon t_1 - \epsilon t)} e^{-i\omega_0 t_1} e^{i\omega_0 t},
\]
where $\Gamma$ is the trace in the spin space, $f_\epsilon(\epsilon)$ is the Fermi distribution, $G^<_{\mathbf{q} \tau}(t, t_1) = \sum_{\mathbf{k} \mathbf{k}'} G_{\mathbf{k} \mathbf{k}' \mathbf{q} \tau}(t, t_1)$, and
\[
G^>_{\mathbf{q} \tau}(t, t_1) = \sum_{\mathbf{k} \mathbf{k}'} G_{\mathbf{k}' \mathbf{k} \mathbf{q} \tau}(t, t_1) = 2 \times 2 \text{ matrices denoting the retarded (advanced) Green’s function and the lesser Green’s function, respectively.}
\]
In the calculation of equation (16), we assume that the dominant contributions to the tunneling stem from the electrons near the Fermi level, and hence the linewidth function is independent of $k$. Thus, we have
\[
\Gamma_\epsilon = \begin{pmatrix} \Gamma^>_\epsilon & 0 \\ 0 & -\Gamma^<_\epsilon \end{pmatrix},
\]
with $\Gamma^>_\epsilon = 2\pi \sum_{\mathbf{q} \mathbf{k} \mathbf{k}'} T_{\mathbf{q} \mathbf{k} \mathbf{k}'} \delta(\epsilon - \epsilon_{\mathbf{k} \tau})$. In order to solve equation (16), we need to calculate the Green’s functions $G^{\tau\tau, >}(t, t')$ and $G^{\tau\tau, <}(t, t')$. Using the equation of motion method, we get
\[
G^{\tau\tau, >}(t, t') = \delta(t-t') G^{\tau\tau, >}(t-t') + \int dt_1 dt_2 G^{\tau\tau, >}(t-t_1) \Sigma^{\tau\tau, >}(t_1, t_2) G^{\tau\tau, >}(t_2, t'),
\]
where
\[
\Sigma^{\tau\tau, >}(t, t') = \int \frac{d\epsilon}{2\pi} \frac{\pi}{\omega_0} e^{-i(\epsilon t' - \epsilon t)} e^{i\omega_0 t} e^{-i\omega_0 t'},
\]
with $\Sigma^0 = -\frac{1}{2} (\Gamma^>_0 + \Gamma^<_0)$, and $G^{\tau\tau, >}(t)$ is the retarded Green’s function of graphene without the coupling of the electrodes, and can be obtained by a straightforward calculation. The detailed expressions are given in the appendix. $G^{\tau\tau, <}(t, t')$ is related to $G^{\tau\tau, >}(t, t')$ through the Keldysh equation
\[
G^{\tau\tau, <}(t, t') = \int dt_1 dt_2 G^{\tau\tau, >}(t_1, t_2) \Sigma^{\tau\tau, <}(t_1, t_2) G^{\tau\tau, >}(t_2, t'),
\]
where
\[
\Sigma^{\tau\tau, <}(t, t') = i \int \frac{d\epsilon}{2\pi} \left[ \Gamma^>_0 f_\epsilon(\epsilon) + \Gamma^<_0 f_\epsilon(\epsilon) \right] \times e^{-i(\epsilon t' - \epsilon t)} e^{i\omega_0 t} e^{-i\omega_0 t'}. \]
Notice that the Green’s functions in equations (18) and (20) do not depend only on the difference of the two time variables, thus one should take a generalized Fourier expansion as [42]
\[
G(t, t') = \frac{1}{2\pi} \sum_\mathbf{n} \int d\epsilon e^{-i\epsilon(t-t')} e^{i\omega_0 \epsilon/2} G(\epsilon, \epsilon + n\omega_0/2).
\]
Hereafter we shall use the simple notation
\[
G_{nm}(\epsilon) = G(\epsilon + n\omega_0/2, \epsilon + n\omega_0/2).
\]
Evidently, the different Fourier components satisfy the relation
\[
G_{nm} = G_{n-m,0}(\epsilon + n\omega_0/2).
\]
From equation (18), the Fourier components of the Green’s function can be expressed as
\[
G_{mm, n}^{\tau\tau, >} = g_{mm, n}^{\tau\tau, >} + \epsilon_m \delta_{mm, n} G_{mm, n}^{\tau\tau, >} + V_{m, m-2} G_{m-2, m}^{\tau\tau, >} + V_{m, m+2} G_{m+2, m}^{\tau\tau, >},
\]
with $\epsilon_m = \delta_{mm, n} \Sigma^{\tau\tau, >}$, and hence
\[
\Sigma^{\tau\tau, >}_m = \delta_{mm, n} \Sigma^{\tau\tau, >}, \epsilon_m = \delta_{mm, n} \Sigma^{\tau\tau, >},
\]
\[
V_{m, m-2} = \begin{pmatrix} g_{m, m-2}^{\tau\tau, >} \delta_{mm, n} \Sigma^{\tau\tau, >} \\ 0 \\ \epsilon_m \delta_{mm, n} \Sigma^{\tau\tau, >} \end{pmatrix},
\]
\[
V_{m, m+2} = \begin{pmatrix} 0 \\ g_{m, m+2}^{\tau\tau, >} \delta_{mm, n} \Sigma^{\tau\tau, >} \\ \epsilon_m \delta_{mm, n} \Sigma^{\tau\tau, >} \end{pmatrix},
\]
with $g_{mm, n}^{\tau\tau, >} = \delta_{mm, n} \tilde{\rho}_{\epsilon_m} + n\omega_0/2)$. Note that equation (23) is formally equivalent to the ones describing the motion of electrons in a tight-binding linear chain with site energies and nearest-neighbor coupling. Its solution can be derived by using the conventional recursive technique [42].

Carrying out the generalized Fourier transformation to equation (20), and substituting it in equation (16) together with equation (23), we finally obtain the time-averaged current (under time average the displacement current for the central region is zero since it is given by the charge
accumulation in the central region over the time period and tends to zero for large time duration [36, 41])

$$\langle I \rangle = \frac{e}{h} \sum_{\tau \tau', n, a} \int \frac{d\epsilon}{2\pi} \left[ G_{0,n,a}^{\tau\tau'}(\epsilon) - G_{0,n,a}^{\tau\tau'}(\epsilon) \right] \Sigma_{\tau,\tau',n}(\epsilon)$$

where

$$G_{\tau,\tau',n,a}^{\tau\tau'}(\epsilon) = \sum_{n,L,R} G_{n,L,R}^{\tau\tau'}(\epsilon) \Sigma_{\tau,\tau',n}(\epsilon) G_{\tau,\tau',n,a}^{\tau\tau'}(\epsilon).$$

In equation (24), we further set the symmetrical voltage division $V_{UL,R} = E_F \pm 1/2eV$, and put $E_F = 0$ in the numerical calculations. The expression (24) is the central result of this paper, and allows one to describe the spin-polarized transport in a nonlinear response regime, i.e., high bias voltage or high strength of the radiation fields. Please note the difference from our previous study [13] in which we studied the influence of a time-dependent chemical potential (realized as an oscillating bias voltage). In the present study the radiation field, coupled to the carrier as a vector potential, induces a dynamical gap in the quasiparticle spectrum. This, in turn, leads to a marked change in the transport properties and TMR, as shown below. The TMR is deduced according to the conventional definition

$$\text{TMR} = \frac{I(0) - I(\pi)}{I(0)}.$$
The reason for this difference is that the sidebands for an oscillating gate voltage are produced due to the modulation of each of the quasiparticle levels by the gate voltage, while those for the present case are caused by the electron scattering between the bands. When the electrons are injected from the electrode to the graphene, the radiation field results in multiple electron scattering events. Each scattering process provides a sideband. When the bias windows cross this sideband, resonant tunneling occurs leading to the appearance of the resonant peak. When the radiation strength is enhanced, the dynamical gap increases. Thus the sidebands fall into the gap at some stage, and the scattering probability of the electrons is suppressed due to the linear density of states (DOS) in graphene which explains the decrease of the peaks. Another pronounced feature is that the positions of the peaks shift in the direction away from the origin with increasing radiation strength, which also differs from the case of a gate voltage [13], where the positions of the peaks are fixed. This is due to the increase of the dynamical gap lifting the photon-assisted sidebands. Figure 1(c) shows the bias dependence of the differential conductance for different frequencies $\omega_0$ in the parallel configuration of the electrodes. With increasing frequency, the dynamical gap diminishes leading to a decrease of the width of the zero value region and to a shift of the peaks towards the origin. In particular, one can find that the amplitudes of the peaks become larger when the frequency grows. The reason is that the frequency increase lifts the photon-induced sidebands, which causes the increase of the DOS for the sidebands, thus enhancing the electron tunneling through the graphene sheet.

Figure 2(a) shows the bias dependence of the differential conductance for the different polarizations $p$ in the parallel electrode configuration. In the differential conductance the peaks are viewed as stemming from the resonant contributions of the sidebands. The other part may be attributed to the off-resonant tunneling. It is observed that the latter contribution is almost independent of $p$, while the former exhibits a sizable dependence on the polarization. This can be understood roughly in that the sidebands become more spin-selective with increasing polarization, thus gradually turning off the spin-flip channel (mixing channel) on the sidebands. This decreases the differential conductance for the resonant states. For very weak peaks, the contribution from the background off-resonant tunneling may compete with those from the resonant states, making the total change vague. Figure 2(b) shows the bias dependence of the differential conductance for different gate voltages. It is found that the oscillation peaks and the zero value region shift towards the positive direction of the bias voltage since the positions of the Dirac points vary due to the gate voltage. The additional interesting phenomenon is that the amplitudes of the oscillation peaks increase with increasing gate voltage. This is because the gate voltage lifts the quasiparticle level of the central graphene in the bias windows; the linear DOS of graphene enhances the tunneling probability from the ferromagnetic electrode to graphene.

The differential conductance as a function of the radiation strength for different radiation frequencies in the parallel configuration of electrode magnetizations is shown in figure 3(a). The differential conductance with varying radiation field strength exhibits several resonant peaks. This can be traced back to the fact that, with a fixed bias window and a fixed frequency, the dynamical gap in the quasiparticle spectrum of graphene becomes larger for stronger radiation field strength, which pushes the sidebands upwards. When the sidebands cross the bias windows, resonant transport occurs, which is then manifested in the conductance peaks. When the radiation field strength is large enough, the differential conductance tends to a constant value which almost does not vary with the frequency, hinting at the conclusion that graphene can be identified as a pure resistance in this situation.

Figure 3(b) shows the dependence of the differential conductance on the radiation frequency $\omega_0$ for different radiation strengths in the parallel configuration of the electrodes. It is clearly seen that when $\omega_0$ approaches zero (note that the magnitude of the field, i.e. $E_0$, is finite), the differential conductance diminishes owing to the large dynamical gap. With increasing frequency, the dynamical gap shrinks, which causes lowering in energy of the photon-induced sidebands. Therefore, this inverse process (to figure 3(a)) causes resonant peaks also in the spectrum in the bias windows. It is found further that the interval between the peaks is not uniform since the radiation frequency changes the dynamical gap in a nonlinear way. Moreover, increasing the radiation field strength enlarges the dynamical gap, and moves the resonant peaks in the differential conductance away from the origin.
Figure 3. The differential conductance $G$ as a function of the radiation field strength for different frequencies (a), and as a function of the frequency for different radiation strengths (b). In both cases we consider the parallel configuration of the electrode magnetizations at $eV = 0.5t_g$. The other parameters are taken the same as those of figure 1.

The bias dependence of the TMR for the different radiation strengths and frequencies is shown in figure 4. The TMR as a function of the bias voltage exhibits an oscillatory behavior. We assign this behavior to the combined effects of the radiation field, the graphene and the spin valve properties. In order to further clarify this point, we plot the bias dependence of the TMR for the different polarizations in figure 5. One can observe that the TMR changes in a nonlinear manner: the TMR values in the off-resonant regions become larger than those at resonant points. This is because graphene remains insulator-like in off-resonant tunneling, the ballistic spin tunneling enhances the TMR. In addition, at high bias voltage, the oscillation peaks decay sharply due to the suppression of the multi-photon processes in graphene. With increasing radiation strength (or frequency), the oscillation is enhanced (or weakened) due to the interplay between the spin valve effect and the dynamical gap (see figure 4).

4. Summary

In conclusion, we studied theoretically the spin-polarized transport through a graphene spin valve device assisted by a linearly polarized, monochromatic laser field and in the presence of a dc bias and a gate voltage. The method is based on the standard Keldysh nonequilibrium Green’s function approach. We find that the bias dependence of the differential conductance exhibits additional peaks due to the resonant transport through the photon-assisted sidebands. The resonant interaction of the quasiparticle in the graphene sheet with the radiation field causes a dynamical gap in the quasiparticle spectrum. When the bias voltage lies inside this gap region, the differential conductance displays a zero value region situated symmetrically around the zero bias. The value of the dynamical gap depends on the strength and the frequency of the external radiation field; thus the width of this zero value region in the differential conductance can be tuned by changing the radiation strength and the frequency, which also causes a shift of the resonant peaks in the resonant transport through the photon-assisted sidebands.
the differential conductance. We explored the behavior of the peaks in the differential conductance with varying radiation field strength and frequency and revealed the dependence on the spin polarization of the ferromagnetic electrodes. We also demonstrated that the combined effects of the radiation field, the graphene and the spin valve properties result in an oscillatory behavior in the TMR. This oscillatory amplitude can be varied by changing the parameters of the radiation field.

Acknowledgments

The work of KHD is supported by DAAD (Germany) and by the National Natural Science Foundation of China (Grant No 10904007), and the construct program of the key discipline in Hunan Province, China. JB and ZGZ are supported by DFG, Germany.

Appendix

In this appendix, we present the analytical results for the Green’s function $g^{\tau,r}_{\mu,s}(t)$ for the graphene without the coupling to the electrodes. From the Hamiltonian $H_G$, we find by using the equation of motion method that

$$
\frac{dg^{\tau,r}_{\mu,s}(t)}{dt} = \int \frac{d\epsilon}{2\pi} g^{\tau,r}_{\mu,s}(\epsilon)e^{-i\omega t},
$$

(A.1)

where $g^{\tau,r}_{\mu,s}(\epsilon) = \sum_k g^{\tau,r}_k(\epsilon)$ with

$$
g^{\tau,r}_k(\epsilon) = \frac{\epsilon - \epsilon_{-k}}{\Delta} \frac{\Delta}{(\epsilon - \epsilon_{+k})(\epsilon - \epsilon_{-k}) - \Delta^2}.
$$

(A.2)

To solve $g^{\tau,r}_k(\epsilon)$, we need to convert the summation over $k$ into a integral in two-dimensional momentum space. After a straightforward calculation, we have, for $|\epsilon| < \Delta$,

$$
g^{\tau,r}_{11} = \frac{1}{2\pi V_F^2} \left| \frac{\epsilon - V_g}{\frac{D - V_g - \omega_0/2}{\omega_0^2/4 + \Delta^2 - (\epsilon - V_g)^2}} \right|

- \sqrt{\Delta^2 - (\epsilon - V_g)^2} \left[ \arctan \frac{D - V_g - \omega_0/2}{\sqrt{\Delta^2 - (\epsilon - V_g)^2}} \right]

- \arctan \frac{-\omega_0/2}{\sqrt{\Delta^2 - (\epsilon - V_g)^2}} + D - V_g \right],
$$

(A.3)

$$
g^{\tau,r}_{22} = \frac{1}{2\pi V_F^2} \left| \frac{\epsilon - V_g}{\frac{D - V_g - \omega_0/2}{\omega_0^2/4 + \Delta^2 - (\epsilon - V_g)^2}} \right|

+ \sqrt{\Delta^2 - (\epsilon - V_g)^2} \left[ \arctan \frac{D - V_g - \omega_0/2}{\sqrt{\Delta^2 - (\epsilon - V_g)^2}} \right]

- \arctan \frac{-\omega_0/2}{\sqrt{\Delta^2 - (\epsilon - V_g)^2}} + D - V_g \right],
$$

(A.4)

For $|\epsilon| > \Delta$,

$$
g^{\tau,r}_{11} = \frac{1}{4\pi V_F^2} \left| \frac{\epsilon - V_g}{\frac{D - V_g - \omega_0/2}{\omega_0^2/4 + \Delta^2 - (\epsilon - V_g)^2}} \right|

- (\epsilon - V_g)^2 + \Delta^2 \left| \frac{D - V_g - \omega_0/2}{\omega_0^2/4 + \Delta^2 - (\epsilon - V_g)^2} \right|^{-1} \right|

- \sqrt{\Delta^2 - (\epsilon - V_g)^2} \left[ \arctan \frac{D - V_g - \omega_0/2}{\sqrt{\Delta^2 - (\epsilon - V_g)^2}} \right]

+ \sqrt{\Delta^2 - (\epsilon - V_g)^2} \left[ \arctan \frac{-\omega_0/2}{\sqrt{\Delta^2 - (\epsilon - V_g)^2}} \right]

\times (\epsilon - V_g + \sqrt{(\epsilon - V_g)^2 - \Delta^2})\right| + 2D - 2V_g

+ \frac{i}{4V_F^2} \text{sgn}(\epsilon - V_g)

\times \left\{ \begin{array}{l}
\theta \left( \Delta < |\epsilon - V_g| < \sqrt{\frac{\omega_0^2}{4} + \Delta^2} \right)

\times (\epsilon - V_g - \sqrt{(\epsilon - V_g)^2 - \Delta^2})

- \theta \left( \Delta < |\epsilon - V_g| < \left( D - V_g - \frac{\omega_0}{2} \right)^2 + \Delta^2 \right)

\times (\epsilon - V_g + \sqrt{(\epsilon - V_g)^2 - \Delta^2})
\end{array} \right\},
$$

(A.6)
\[ D \] is the high-energy cutoff of the graphene bandwidth.

References

[1] Duan X, Huang Y, Cui Y, Wang J and Lieber C M 2001 Nature 409 66
[2] Marcus M S, Simmons J M, Castellini O M, Hammers R J and Eriksson M A 2006 J. Appl. Phys. 100 084306
[3] Freitag M, Martin Y, Misewicj J A, Martel R and Avouris P 2003 Nano Lett. 3 1067
[4] Ohno Y, Kishimoto S, Mizutani T, Okazaki T and Shinozaka H 2004 Appl. Phys. Lett. 84 1368
[5] Wang J, Gudiksen M S, Duan X, Cui Y and Lieber C M 2001 Science 293 1455
[6] Novoselov K S, Geim A K, Morozov S V, Jiang D, Katsnelson M I, Grigorieva I V, Dubonos S V and Firsov A 2005 Nature 438 197
[7] Katsnelson M I, Novoselov K S and Geim A K 2006 Nature Phys. 2 620
[8] Geim A K and Novoselov K S 2007 Nature Mater. 6 183
[9] Peres N M R, Guinea F and Castro Neto A H 2006 Phys. Rev. B 73 125411
[10] Falkovsky L A and Varlamov A A 2007 Eur. Phys. J. B 56 281
[11] Gusynin V P and Sharapov S O 2006 Phys. Rev. B 73 245411
[12] Trauzettel B, Blanter Y M and Morpurgo A F 2007 Phys. Rev. B 75 155405
[13] Ding K H, Zhu Z-G and Berakdar J 2011 Phys. Rev. B 84 115433
[14] Gusynin V P, Sharapov S G and Carbotte J P 2007 Phys. Rev. B 75 155407
[15] Falkovsky L A and Pershoguba S S 2007 Phys. Rev. B 76 155410
[16] Abergel D S L and Falco V I 2007 Phys. Rev. B 75 155430
[17] Huang E H and Das Sarma S 2007 Phys. Rev. B 75 205418
[18] Wunsch B, Stauber T, Sols F and Guinea F 2006 New J. Phys. 8 318
[19] Apalkov V, Wang X-F and Chakraborty T 2007 Int. J. Mod. Phys. B 21 1165
[20] Vafek O 2006 Phys. Rev. Lett. 97 266406
[21] Mikhailov S A and Ziegler K 2007 Phys. Rev. Lett. 99 016803
[22] Mikhailov S A and Ziegler K 2008 J. Phys.: Condens. Matter 20 384204
[23] Moskalenko A S and Berakdar J 2009 Phys. Rev. B 80 193407
[24] Sensale-Rodriguez B, Fang T, Yan R, Kelly M M, Jena D, Liu L and Xing H 2011 Appl. Phys. Lett. 99 113104
[25] Fusilier S V, Fusilier M V and Efetov K B 2008 Phys. Rev. B 77 045407
[26] Hill E W, Geim A K, Novoselov K, Schedin F and Black P 2006 IEEE Trans. Magn. 42 2694
[27] Tomblin N, Joza C, Popincute M, Jonkmans H T and Van Wees B J 2007 Nature 448 571
[28] Ohishi M, Shiraishi M, Nozaki T, Shinjo T and Suzuki Y 2007 Japan. J. Appl. Phys. 46 L605
[29] Cho S, Chen Y-F and Fuhrer M S 2007 Appl. Phys. Lett. 91 123105
[30] Wang W H, Pi K, Li Y, Chiang Y F, Wei P, Shi J and Kawakami R K 2008 Phys. Rev. B 77 020402(R)
[31] Ding K H, Zhu Z-G and Berakdar J 2009 Phys. Rev. B 79 045405
[32] Ding K H, Zhu Z-G and Berakdar J 2009 Europhys. Lett. 88 58001
[33] Ding K H, Zhu Z-G and Berakdar J 2008 J. Phys.: Condens. Matter 20 345228
[34] Maassen J, Ji W and Guo H 2011 Nano Lett. 11 151
[35] Chen J-C, Cheng S-G, Shen S-Q and Sun Q-F 2010 J. Phys.: Condens. Matter 22 035301
[36] Ding K H, Zhu Z-G and Berakdar J 2010 Phys. Rev. B 82 155143
[37] Haug H and Jauho A P 1998 Quantum Kinetics in Transport and Optics of Semiconductors (Berlin: Springer)
[38] Rammer J 2004 Quantum Transport Theory (Boulder, CO: Westview Press)
[39] Castro Neto A H, Guinea F, Peres N M R, Novoselov K S and Geim A K 2009 Rev. Mod. Phys. 81 109
[40] Gerry C C and Knight P L 2005 Introductory Quantum Optics (Cambridge: Cambridge University Press)
[41] Zhang X, Wang J and Zhang S C 2010 Phys. Rev. B 82 245107
[42] Zhu Z-G, Su G, Zheng Q-R and Jin B 2004 Phys. Rev. B 70 174403
[43] Levy Yeyati A and Flores F 1991 Phys. Rev. B 44 9020