Valley controlled spin-transfer torque in ferromagnetic graphene junctions

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Abstract
The presence of the valley degree of freedom in graphene leads to the valleytronics, in which information is encoded by the valley quantum number of the electron. We propose a valley controlled spin-transfer torque (STT) in graphene-based normal/normal/ferromagnetic junctions with the normal lead irradiated by the off-resonant circularly polarized light. The interplay of the spin–orbit interaction and the staggered potential in the central normal part results in the coupling between the valley and spin degrees of freedom, so a valley dependent spin polarized current can be demonstrated, which can exert a valley controlled STT on the ferromagnetic lead. The amplitude of the STT can be manipulated by the intensity of the light, the Fermi energy and the magnetization direction of the ferromagnetic lead. This valley controlled STT may find potential application in future valleytronics and spintronics.

1. Introduction

Over the past decade, motivated by the development of spintronics, the presence of the valley degree of freedom in graphene leads to the valleytronics [1, 2], whose goal is to manipulate the valley degree of freedom and search for its potential applications in semiconductor technologies and quantum information. Much substantial progress in valleytronics has been made recently, such as quantum valley Hall effect [3–6], valley polarization controlled by circularly polarized light in molybdenum disulfide [7], and valley and spin currents in silicene junctions [8–10]. Interestingly, if the intrinsic spin–orbit interaction and the staggered potential coexist in graphene-like materials, the band structure becomes spin-valley coupling, so one can control the spin polarized current by the valley degree of freedom [8–11].

Recently, on the other hand, the spin-transfer torque (STT) has also attracted much attention [12–19]. When a spin polarized current is injected into a ferromagnetic layer with a magnetization misaligned to the spin polarization of the current, it can transfer spin angular momentum to the ferromagnetic layer, and hence exerts a torque on the magnetic moments of the ferromagnetic layer, which may change the magnetization orientation of the ferromagnetic layer. The STT effect has the potential application in random access memory, which offers an all-electrical read and write process. The STT was theoretically predicted by Slonczewski [12] and Berger [13], and then it has been extensively confirmed experimentally [14]. Recently, the STT in ferromagnetic graphene junctions has also been reported. Yokoyama and Linder [15] demonstrated that both the magnitude and the sign of the STT can be controlled by means of the gate voltage in a bulk ferromagnetic/normal/ferromagnetic graphene junction. Then Ding et al [16] investigated theoretically the effect of strain on the STT in a zigzag-edged graphene nanoribbon spin–valve device. Later Zhang et al [17] studied the helical spin polarized current induced STT in graphene-based normal/topological insulator/ferromagnetic junctions. Although the STT in ferromagnetic graphene junctions has been investigated in [15–17], the effect of the valley degree of freedom on the STT is not explored. If the coupling between the valley and spin degrees of freedom exists, a valley controlled spin polarized current is generated. Naturally, it is interesting to discuss the effect of the valley degree of freedom on the STT in a ferromagnetic graphene junction.
In this work, we predict a valley controlled STT in graphene-based normal/normal/ferromagnetic (N1/N2/F) junctions, where the N1 is irradiated by the off-resonant circularly polarized light. In the N1, the valley polarization can be modulated by the interaction between the staggered potential and the light. While in the N2, the interplay of the spin–orbit interaction and the staggered potential leads to the coupling between the valley and spin degrees of freedom. In this case one can control the valley dependent spin polarized current by the intensity of the light, the Fermi level and Rashba spin–orbit interaction on the STT are discussed.

The rest of this paper is organized as follows. In section 2 we give the Hamiltonian of the N1/N2/F junctions, discretize the Hamiltonian in the basis \( |x_i\rangle \otimes |k_j\rangle \), and then present the formula of the valley dependent STT by the Keldysh non-equilibrium Green’s function method. In section 3, numerical results and detailed discussions are demonstrated. Finally, in section 4 we summarize the main conclusions of this work.

2. Model and formulation

We consider a graphene-based two dimensional N1/N2/F junction (see figure 1 (a)) with the interfaces located at \( x = 0 \) and \( x = L \), where \( L \) is the width of the N2. The N1 \( (x < 0) \) is irradiated by off-resonant circularly polarized light that requires \( h\omega \gg \gamma \) in principle, where \( \omega \) is the frequency of light and \( \gamma \) is the nearest-neighboring hopping energy in graphene [20–24]. In the central part N2 the spin–orbit interaction and the staggered potential are considered, which leads to the coupling between the valley and spin degrees of freedom. The ferromagnetic electrode deposited on top of the graphene sheet \( (x > L) \) induces a finite exchange field \( \mathbf{h} = h(\cos \theta, 0, \sin \theta) \) [8, 15, 25], where \( h \) is the magnitude of exchange field and \( \theta \) describes the direction angle of the magnetization. In the low-energy approximation, the Hamiltonian of the present junctions is expressed as [8, 9, 21, 26]

\[
H^{\text{eff}} = \sigma_0 \otimes (\epsilon N_F \sigma_{x} \tau_{z} + k_z \tau_{y}) + (\eta \lambda_{\sigma} - \lambda_{c}) \tau_{z} + \eta \lambda_{\sigma} \sigma_{x} \otimes \tau_{z} - \sigma \cdot \mathbf{h} \otimes \tau_{b}.
\]

Here \( \eta = +1 (-1) \) represents the \( K (K') \) valley with \( v_F \) the velocity of electrons. \( \tau_{j} \) and \( \sigma_{j} (j = x, y, z, 0) \) are the Pauli matrices and unit matrices in valley and spin space, respectively. \( \eta \lambda_{\sigma} \sigma_{x} \otimes \tau_{z} \) is intrinsic spin–orbit coupling term in graphene. As indicated in Kane and Mele investigation [26], graphene will be driven into topological phase when this term in graphene is enhanced. In general, the spin–orbit interaction is very weak in graphene. However, many works suggested that spin–orbit interaction can be enhanced by the substrates or adatom deposition [27–30], which results in the nontrivial topological phase. \( \lambda_{c} \) is related to the intensity of the off-resonant circularly polarized light and provides an additional site energy of a sublattice in the N1. By using the Floquet theory [20–24] \( \lambda_{c} \) can be written as \( \lambda_{c} = \frac{8\pi \sigma \xi v_F^2}{\omega^2} \), where \( \xi = (+) \) corresponds to the right-hand (left-hand) circularly polarized light with the frequency \( \omega \), the intensity \( I \), and the fine structure constant

Figure 1. (a) Schematic plot of a N1/N2/F junction, where a graphene monolayer sheet in the xy plane is grown on the substrates. Here the staggered potential induced by the substrates is assumed to be finite in the N1 and N2 but zero in the F. (b) The band structure of the N1/N2/F junctions. In the N1, the spin is degenerate, and due to the interaction between the staggered potential (SP) and the light, a valley-dependent band gap exists. In the N2, the interplay of the spin–orbit interaction (SOC) and the SP leads to the coupling between the valley and spin degrees of freedom. While in the F the valley is degenerate.
Before presenting the numerical results for the valley dependent STT in the N1 channel, we need to discuss some details about the system. The valley dependent STT is studied by the non-equilibrium Green’s function method. The retard Green’s function of the central part N2 can be calculated by the following expression:

\[ G'(\varepsilon) = \left[ (\varepsilon + i\eta)I - H_C - \sum_{L} - \sum_{R} \right]^{-1}, \]

where the self-energy is \( \Sigma^{R}_{L,R} = H_{R,L} \delta_{R,L} - iE_0 \delta_{R,L} \delta_{R+1,L+1} + iE_0 \delta_{R,L} \delta_{R-1,L-1} \), and \( H_{R,L} \) is the coupling matrix between the left (right) lead and the central part, and \( H_C \) is the Hamiltonian of the N2. For a small bias voltage, the valley dependent STT per unit of the bias voltage in zero temperature can be obtained as

\[ \tau^{Rk}_{x} = \frac{\varepsilon}{4\pi} \sum_{k} Tr \left\{ (G'(\mu_L)\Gamma_L(\mu_L)G'^{-1}(\mu_L)\Gamma_R(\mu_R))(\sigma_x \cos \theta - \sigma_y \sin \theta) \right\}, \]

where \( \mu_L \) and \( \mu_R \) are the Fermi level. By using equation (6), we can obtain the total STT \( \tau^{R^2} = \tau^{R}_{x} + \tau^{R}_{z} \).

### 3. Results and discussions

Before presenting the numerical results for the valley dependent STT in the N1/N2/F junctions, we first make a physical analysis for the valley dependent STT. As shown in equation (2), in the N1, the spin is degenerate and the valley polarization can be manipulated by stagger potential and the off-resonant circularly polarized light. While from equation (1) the dispersion relation of the N2 can be written as \( \varepsilon_{n} = n\sqrt{-(\lambda_0 - \eta\lambda_\omega)^2 + (\hbar v_F k)^2} \), where \( n = \pm \) represents the conduction (valence) band, and \( k = \sqrt{k_x^2 + k_y^2} \) is the modulus of wave vector \( k \). It is noted that for the electrons in the valley where there exists an energy gap \( \Delta E \), which \( \pm \) can be tuned by the parameters \( \lambda_0 \) and \( \lambda_\omega \). The shift of the incident electron should be satisfy \( |\varepsilon| > \Delta E/2 \) to generate propagating incident modes in the valley of the normal lead.

Due to the translation invariant along the y-axis, the transversal wave vector \( k_y \) of the incident electron must be conserved. Following the [34], we can discretize the Hamiltonian (1) in the basis \{\( |x_i\rangle \otimes |k_j\rangle \} \)

\[ H^{0}_{i,i'} = H_{i,i'} \delta_{i,i'} - \varepsilon \delta_{i,i'+1} + iE_0 \delta_{i,i'} \delta_{i,i'+1} + \varepsilon \delta_{i,i'} \delta_{i,i'+1}, \]

where \( E_0 = \frac{\hbar v_F a}{a} \) and \( a = x_{i+1} - x_i \) with \( a = 0.5 \) nm [35] is the mesh spacing along the x direction (\( i \leq N \) for the left lead; \( 1 \leq i \leq N \) for the central part, and \( i \geq N + 1 \) for the right lead). \( H_{i,i'} \) is given as

\[ H_{i,i'} = \delta_{i,i'} \left[ \right( \varepsilon + i\eta) \delta_{i,i'} - \sum_{L} - \sum_{R} \right]^{-1}, \]

where the self-energy is \( \Sigma^{R}_{L,R} = H_{R,L} \delta_{R,L} \), \( H_{R,L} \) is the surface retarded Green’s function of the left (right) lead, and \( H_C \) is the Hamiltonian of the N2. For a small bias voltage, the valley dependent STT per unit of the bias voltage in zero temperature can be obtained as

\[ \tau^{Rk}_{x} = \frac{\varepsilon}{4\pi} \sum_{k} Tr \left\{ (G'(\mu_L)\Gamma_L(\mu_L)G'^{-1}(\mu_L)\Gamma_R(\mu_R))(\sigma_x \cos \theta - \sigma_y \sin \theta) \right\}, \]

where \( \mu_L \) is the Fermi level. By using equation (6), we can obtain the total STT \( \tau^{R^2} = \tau^{R}_{x} + \tau^{R}_{z} \).
amplitude of $\tau_{Rx}$ first increases and then begins to decrease, so in order to get a large $\tau_{Rx}$, one should choose an appropriate parameter $\lambda_\omega$. We can understand this behavior as follows. For a finite positive $\lambda_\omega$, because the density of states of the incident electrons in the $K$ valley nonmonotonically depends on $\lambda_\omega$, the current coming from the $K$ valley first increases and then decreases with $\lambda_\omega$, leading to a nonmonotonic dependence of $K_{Rx}$ on $\lambda_\omega$. While for the parameters taken here the density of states of the incident electrons in the $K'$ valley monotonically decreases with $\lambda_\omega$ and becomes zero for large $\lambda_\omega$, so $K_{Rx}'$ decreases and becomes zero for large $\lambda_\omega$.

As discussed above, the off-resonant circularly polarized light strongly influences on the valley dependent STT, so it is interesting to analyze the effect of $\lambda_\omega$ on the valley dependent STT in detail. Figure 3(a) displays the valley dependent STT as a function of $\theta$ at $\omega = \pi/4$ (black lines), $\pi/2$ (read lines) and $3\pi/4$ (blue lines). As shown in equation (2), the band structure of the $\eta$ valley in the N1 has an band gap $E_g^\eta$, which can be tuned by $\lambda_\omega$.

For the $K$ valley $\mu_F > E_g^K/2$ should be satisfied to generate propagating incident modes in the N1, so for the parameters $\mu_F = 30$ meV and $\lambda_\omega = 20$ meV taken here, $\tau_{Rx}'$ nonmonotonically depends on $\lambda_\omega$ in the regime of $-10$ meV $< \lambda_\omega < 50$ meV. On the other hand, for the $K'$ valley $\mu_F > E_g^{K'}/2$ is needed to generate propagating incident modes, thus $\tau_{Rx}'$ nonmonotonically depends on $\lambda_\omega$ in the regime of $-50$ meV $< \lambda_\omega < 10$ meV. In this case the curves of $\tau_{Rx}$ versus $\lambda_\omega$ can divide into three regimes. (1) For $-50$ meV $< \lambda_\omega < -10$ meV $\tau_{Rx}$ is positive and only originates from $\tau_{Rx}'$. (2) For $-10$ meV $< \lambda_\omega < 10$ meV both $\tau_{Rx}$ and $\tau_{Rx}'$ contribute to $\tau_{Rx}$, which changes from positive to negative with increase of $\lambda_\omega$. (3) For $\lambda_\omega > 10$ meV $\tau_{Rx}$ is negative and only comes from $\tau_{Rx}'. At \theta = \pi/2$, due to the presence of $\tau_{Rx}'(\lambda_\omega) = \tau_{Rx}'(-\lambda_\omega)$, $\tau_{Rx}'$ is an odd function of $\lambda_\omega$ (red dotted line in figure 3(a)). On the other hand, for the other $\theta$, the symmetry of $\tau_{Rx}'(\lambda_\omega) = \tau_{Rx}'(-\lambda_\omega)$ is broken, $\tau_{Rx}'$ is not an odd function of $\lambda_\omega$ any more. In order to explain the behavior of $\tau_{Rx}'$, we divide $\tau_{Rx}'$ into two parts:
Figure 3. (a) $\tau_{\pi R}^E$ (solid lines), $\tau_{\pi L}^E$ (dashed lines) and $\tau_{\pi E}^E$ (dotted lines) versus $\lambda_{\omega}$ at different $\theta$. (b) $-G_z \sin(\phi)$ and (c) $G_x \cos(\phi)$ versus $\lambda_{\omega}$ at different $\theta$. Other parameters are the same as those in figure 2.
\[ \tau_{Q1}^{Rx} = -G_2 \sin \theta \quad \tau_{Q2}^{Rx} = G_\lambda \cos \theta \]

with \( G_\pm = \frac{1}{2} \sum_{\sigma \in \{+,-\}} \mathrm{Tr} (G^\dagger \Gamma_{Q1} \Gamma_1 \sigma_{z(\pm)}) \). For \( \theta = \pi / 2 \) \( \tau_{Q2}^{Rx} \) is zero and \( \tau_{Q1}^{Rx} \) only comes from \( \tau_{Q1}^{Rx} \), which has the symmetry of \( \tau_{Q1}^{Rx}(-\lambda_\omega) = \tau_{Q1}^{Rx}(\lambda_\omega) \), thus \( \tau_{Q1}^{Rx} \) is an odd function of \( \lambda_\omega \).

However, for \( \theta = \pi / 4 \) or \( 3\pi / 4 \) both \( \tau_{Q1}^{Rx} \) and \( \tau_{Q2}^{Rx} \) (\( \tau_{K1}^{Rx} \) and \( \tau_{K2}^{Rx} \)) are finite and contribute to \( \tau_{Q}^{Rx} \) (\( \tau_{K1}^{Rx} \) and \( \tau_{K2}^{Rx} \) have the opposite sign (figure 3(b)), but \( \tau_{K1}^{Rx} \) and \( \tau_{K2}^{Rx} \) have the same sign (figure 3(c))). Because \( \tau_{Q1}^{Rx} \) comes from the combined contribution of \( \tau_{Q1}^{Rx} \) and \( \tau_{Q2}^{Rx} \), which breaks the symmetry of \( \tau_{Q}^{Rx}(\lambda_\omega) = \tau_{Q}^{Rx}(-\lambda_\omega) \), \( \tau_{Q}^{Rx} \) is not an odd function of \( \lambda_\omega \) any more.

Since the STT strongly depends on the magnetization direction of the F (figure 2) and the off-resonant circularly polarized light (figure 3), the STT as a function of \( \theta \) and \( \lambda_\omega \) is plotted in figure 4. It is found that the STT is zero for \( \theta = 0, \pi, 2\pi \), where the spin polarization direction of the incident current is parallel or antiparallel to the magnetization direction of the F. For \( \theta \) near the angles \( \theta = \frac{\pi}{2}, \frac{3\pi}{2} \) the STT can reach maximum. When \( \theta \) is fixed, as discussed in figure 2, with increase of \( \lambda_\omega \), the amplitude of \( \tau_{Q1}^{Rx} \) first increases and then begins to decrease, so in order to get a large \( \tau_{Q1}^{Rx} \) one should choose an appropriate parameter \( \lambda_\omega \). We also find by changing the sign of \( \lambda_\omega \) the STT can convert from the negative to positive, so we can control the STT direction by the handedness of the light.

Furthermore, let us discuss the effect of the Fermi level \( \mu_F \) on the valley dependent STT at \( \lambda_\omega = 10 \) meV (solid lines) and \(-10 \) meV (dashed lines) in figure 5. At \( \lambda_\omega = 10 \) meV \( \tau_{Q1}^{Rx} \) is finite and negative for \( \mu_F > 10 \) meV. With increase of \( \mu_F \) \( \tau_{Q1}^{Rx} \) (blue solid line) first increases with its slope changing at \( \mu_F = 15 \) and
20 meV, and then decreases for $\mu_F > 30$ meV. This can be explained by $\tau_{Kz}^{Rx}$ (black solid line) and $\tau_{Kx}^{Rx}$ (red solid line). When $\mu_F$ is smaller than 10 meV, because $\mu_F$ lies in the band gap of the $K$ valley in the $N_1$, the current is zero, leading to vanishing $\tau_{Kz}^{Rx}$. For 10 meV < $\mu_F$ < 20 meV $\mu_F$ locates at the conduction band of the $K$ valley in the $N_1$, the current becomes finite. We find a finite $\tau_{Kz}^{Rx}$, whose amplitude first increases and then becomes almost constant for 15 meV < $\mu_F$ < 20 meV. This is because in this regime $\tau_{Kz}^{Rx}$ depends on the density of states of the band structure $E_{K_+} = \pm (\hbar v_K k + h)$ and $E_{K_-} = \pm (\hbar v_K k - h)$ in the right $F$, which, respectively, decreases and increases with $\mu_F$, so $\tau_{Kz}^{Rx}$ first increases and then becomes almost constant for 15 meV < $\mu_F$ < 20 meV. For $\mu_F > 20$ meV the density of states of $E_{K_+}$ ($E_{K_-}$) increases with $\mu_F$, so the spin polarized current exerted on the $F$ increases, leading to the increase of $\tau_{Kz}^{Rx}$. Once $\mu_F$ is larger than 30 meV, $\mu_F$ also locates at the conduction band of the $K'$ valley in the $N_1$, $\tau_{Kx}^{Rx}$ increases with $\mu_F$. Because $\tau_{Kz}^{Rx}$ and $\tau_{Kx}^{Rx}$ have opposite sign, $\tau_{Kz}^{Rx}$ relies on the competition of $\tau_{Kz}^{Rx}$ and $\tau_{Kx}^{Rx}$. $\tau_{Kx}^{Rx}$ increases faster than $\tau_{Kz}^{Rx}$, so with further increasing of $\mu_F$, the amplitude of $\tau_{Kz}$ decreases. Due to the presence of the relationship of $\tau_{Kz}^{Rx}(\lambda_\omega)\big|_{\lambda_\omega=10}$ meV, one observes the STTs at $\lambda_\omega = 10$ meV and $\lambda_\omega = 10$ meV have the same amplitude but opposite sign, thus the sign of the STT can be controlled by handedness of the light.

In addition, due to the presence of the structure inversion asymmetry in the $z$ direction, the Rashba spin–orbit interaction $\lambda(\sigma_z \otimes \tau_y - \eta \sigma_y \otimes \tau_z)$ may appear in the $N_2$. In figure 6 we study the effect of the Rashba spin–orbit interaction on $\tau^{Rx}$ at $\theta = \pi / 2$. As shown in figure 6 $\tau^{Rx}$ is an odd function of $\lambda_\omega$ and decreases with $\lambda$. This is because without $\lambda$ the current can arrive at the $N_2/F$ interface with the spin polarization direction along $z$ (for the $K$ valley) or $-z$ (for the $K'$ valley) axis, which is perpendicular to the magnetization direction of the $F$. However, when $\lambda$ is finite, the electrons precess in the process of traveling across the $N_2$, so the spin polarization direction deviates from the $z$ axis, which results in the decrease of $\tau^{Rx}$ with $\lambda$. Therefore in order to obtain a large $\tau^{Rx}$, a small $\lambda$ is needed.

Last we will comment on the experimental feasibility of our results. The staggered potential can be induced by the substrate. Actually, in experiment, the gap induced by the SiC substrate can range from several meV to almost constant for zero, leading to vanishing $\tau^{Rx}$. It should be pointed out that the amplitude of $\tau^{Rx}$ obtained here is comparable to the STT in a previous work, where a ferromagnetic/normal/ferromagnetic junction is
investigated [15]. However, unlike [15], the spin polarized currents in this work originates from the coupling between the valley and spin degrees of freedom, so the STT reported here does not require additional ferromagnetic layer with fixed magnetization.

4. Summary

In summary, we study the valley dependent STT in N$_1$/N$_2$/F junctions. The N$_1$ is irradiated by the off-resonant circularly polarized light, and the valley polarization can be modulated by the interaction between the staggered potential and the light. While in the N$_2$, due to the interplay of the spin–orbit interaction and the staggered potential induced by the substrate, the band structure is spin-valley coupling, so one can control the valley dependent spin polarized current by the valley degree of freedom, which exerts a valley controlled STT on the F. The effects of the intensity of the light, the Fermi level and Rashba spin–orbit interaction on the STT are investigated. The valley controlled STT reported here suggests the ferromagnetic graphene junction ideal for very efficient magnetization manipulation of magnetic materials without external magnetic fields.

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