Controllable spin transport in ferromagnetic graphene junctions

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We study spin transport in normal/ferromagnetic/normal graphene junctions where a gate electrode is attached to the ferromagnetic graphene. We find that, due to the exchange field of the ferromagnetic graphene, spin current through the junctions has an oscillatory behavior with respect to the chemical potential in the ferromagnetic graphene, which can be tuned by the gate voltage. In particular, we obtain a spin current reversal controllable by the gate voltage. Our prediction of high controllability of spin transport in ferromagnetic graphene junctions may contribute to the development of spintronics.

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There is rapidly growing attention to graphene because it has a rich potential from both fundamental and applied physics points of view.[1–3] Graphene has a two-dimensional honeycomb network of carbon atoms and, as a result, electrons in graphene are governed by the Dirac equation, which provides a bridge between condensed matter physics and quantum electrodynamics, thus fascinating theoreticians working toward the ultimate goal of a "unified" theory.[4,5] The progress of practical fabrication techniques for single graphene sheets has allowed experimental study of this system, which has attracted tremendous interest from the scientific community.[6–8] There have been intensive studies on graphene to this date, for instance, half-integer and unconventional quantum Hall effect,[7,9,10] observation of minimum conductivity[8] and a bipolar supercurrent.[11]

Graphene has many important properties for applications: it exhibits gate-voltage-controlled carrier conduction, high field-effect mobilities, and a small spin-orbit interaction.[12,13] Thus, it is extremely promising for future technologies, especially spintronics. Motivated by this, some studies on graphene have been implemented. Spin injection into a graphene thin film has been successfully demonstrated by using nonlocal magnetoresistance measurements.[14–16] The possibility of inducing ferromagnetic correlations in graphene due to the proximity effect by magnetic gates in close proximity to graphene has also been discussed.[17] In these works, induced ferromagnetism in graphene is "extrinsic."

On the other hand, there is the attempt to induce ferromagnetism "intrinsically" in graphene. It has recently been predicted that zigzag-edge graphene nanoribbons become half metallic on exposure to an external transverse electric field due to the different chemical potential shift at the edges,[18–20] which indicates the high controllability of ferromagnetism in graphene and hence opens the possibility of spintronics applications for graphene. In view of this, the study of spin transport in ferromagnetic graphene is very timely and desirable for the development of spintronics.

Stimulated by this, in this paper, we study spin transport in normal/ferromagnetic/normal graphene junctions where a gate electrode is attached to the ferromagnetic graphene. We find that, due to the exchange field in the ferromagnetic graphene, spin current through these systems has an oscillatory behavior with respect to the chemical potential in the ferromagnetic graphene, which can be tuned by the gate electrode.[7–9] In particular, we find a spin current reversal controllable by the gate voltage. Since the exchange field of graphene is also tunable by an in-plane external electric field and half metallicity can even be induced,[18–20] our prediction of high controllability of spin transport in ferromagnetic graphene junctions may facilitate the development of spintronics. Note that the model itself in this paper is the same as that in Ref. 17 and here we mainly focus on the case with the exchange field comparable to the Fermi energy in the model of Ref. 17.

Now let us explain the formulation. The fermions around the Fermi level in graphene obey a massless relativistic Dirac equation. The Hamiltonian is given by

$$H_{\pm} = v_F(\sigma_x k_x \pm \sigma_y k_y) \tag{1}$$

with Pauli matrices σ_x and σ_y, and the velocity $v_F \approx 10^6$ m/s in graphene. This is roughly 100 times larger than in a normal metal, and thus it is safe to neglect the Coulomb interaction compared to the kinetic energy in graphene.[21,22] The Pauli matrices operate on the two-triangular-sublattice space of the honeycomb structure, corresponding to the A and B atoms. The $\pm$ sign refers to the two so-called valleys of the K and K' points in the Brillouin zone. Also, there is a valley degeneracy, which allows one to consider one of the $H_{\pm}$ set.[23] The linear dispersion relation is a reasonable approximation even for Fermi levels as high as 1 eV,[24] so that the fermions in graphene behave like massless Dirac fermions in the low-energy regime.

We consider a two-dimensional normal/ferromagnetic/normal graphene junction where a gate electrode is attached to the ferromagnetic graphene. This junction may be realized by applying an external transverse electric field to a part of a graphene nanoribbon to make it partially ferromagnetic. See Fig. 1 for a schematic of the model with corresponding dispersion relations. The interfaces are parallel to the y axis and located at $x=0$ and $x=L$. Since there is a valley degeneracy, we focus on the Hamiltonian H_+ with $H_+ = v_F(\sigma_x k_x + \sigma_y k_y) - V(x)$; $V(x) = E_F$ in normal graphenes and $V(x) = E_F + U \pm H$ in the ferromagnetic graphene. Here,

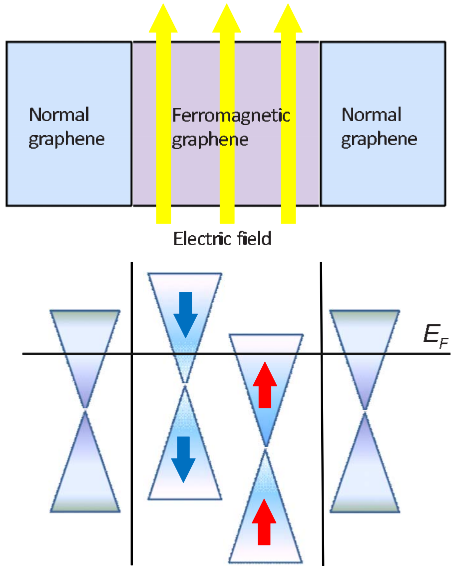

FIG. 1. (Color online) Schematic of the model of a normal/ferromagnetic/normal graphene junction with corresponding Dirac cones. In the ferromagnetic graphene, different Dirac cones correspond to the different chemical potentials of majority and minority spins.

$E_F=v_F k_F$ is the Fermi energy, U is the chemical potential shift tunable by the gate voltage, and H is the exchange field. The $\pm$ signs correspond to majority and minority spins. The wave functions are given by

$$\psi_1 = \begin{pmatrix} 1 \\ e^{i\theta} \end{pmatrix} e^{ip\cos\theta x + ip_y y} + a_\pm \begin{pmatrix} 1 \\ -e^{-i\theta} \end{pmatrix} e^{-ip\cos\theta x + ip_y y}, \quad (2)$$

$$\psi_2 = b_\pm \begin{pmatrix} 1 \\ e^{i\theta'} \end{pmatrix} e^{ip'_\pm \cos\theta' x + ip_y y} + c_\pm \begin{pmatrix} 1 \\ -e^{-i\theta'} \end{pmatrix} e^{-ip'_\pm \cos\theta' x + ip_y y}, \quad (3)$$

$$\psi_3 = d_\pm \begin{pmatrix} 1 \\ e^{i\theta} \end{pmatrix} e^{ip\cos\theta x + ip_y y} \quad (4)$$

with angles of incidence θ and θ', $p=(E+E_F)/v_F$, and $p'_\pm=(E+E_F+U\pm H)/v_F$. Here, ψ_1 and ψ_3 are wave functions in the left and right normal graphenes, respectively, while ψ_2 is a wave function in the ferromagnetic graphene. Because of the translational symmetry in the y direction, the momentum parallel to the y axis is conserved: $p_y=p\sin\theta=p'\sin\theta'$.

By matching the wave functions at the interfaces ($\psi_1=\psi_2$ at $x=0$ and $\psi_2=\psi_3$ at $x=L$), we obtain the coefficients in the wave functions. Note that these conditions are equivalent to $\hat{v}_x\psi_1=\hat{v}_x\psi_2$ at $x=0$ and $\hat{v}_x\psi_2=\hat{v}_x\psi_3$ at $x=L$, with velocity operator $\hat{v}_x=\partial H_+/\partial k_x=v_F\sigma_x$, and hence the current is conserved at the interfaces. The transmission coefficient has the form

$$d_\pm = \frac{\cos\theta\cos\theta' e^{-ipL\cos\theta}}{\cos(p'_\pm L\cos\theta')\cos\theta\cos\theta' - i\sin(p'_\pm L\cos\theta')(1-\sin\theta\sin\theta')}. \quad (5)$$

Then the dimensionless spin-resolved conductances $G_{\uparrow,\downarrow}$ are given by

$$G_{\uparrow,\downarrow} = \frac{1}{2}\int_{-\pi/2}^{\pi/2} d\theta \cos\theta T_{\uparrow,\downarrow}(\theta) \quad (6)$$

with $T_{\uparrow,\downarrow}(\theta)=|d_\pm(\theta)|^2$. Finally, the spin conductance G_s is defined as $G_s=G_\uparrow - G_\downarrow$. Below, we focus on the zero-voltage conductances, namely, we set $E=0$.

Now, let us explain the underlying mechanism of spin manipulation by the gate voltage. In the limit of $|U\pm H| \gg E_F$, we have $\theta' \to 0$ and hence the transmission coefficient of the form

$$d_\pm \to \frac{\cos\theta e^{-ipL\cos\theta}}{\cos\chi_\pm \cos\theta - i\sin\chi_\pm} \quad (7)$$

with $\chi_\pm=\chi\pm\chi_H$, $\chi=UL/v_F$, and $\chi_H=HL/v_F$. The resulting transmission probability is represented as[25]

$$T_{\uparrow,\downarrow}(\theta) \to \frac{\cos^2\theta}{1-\sin^2\theta\cos^2\chi_\pm}. \quad (8)$$

From this expression, we see the π periodicity with respect to $\chi_\pm$ or χ.[17,25–27] We also find that $G_{\uparrow,\downarrow}$ has a maximum (minimum) value of 1 (2/3) at $\chi_\pm=0$ ($\pi/2$). The phase difference between $G_\uparrow$ and $G_\downarrow$ is given by $\chi_+ - \chi_- = 2\chi_H = 2HL/v_F$. If this is equal to the half period $\pi/2$ (that is, $H/E_F=\pi k_F L/4$), we expect a large spin current which oscillates with χ, namely, the gate voltage, because when one of $G_\uparrow$ and $G_\downarrow$ has a maximum at a certain χ, the other has a minimum at the same χ. In this case, the value of G_s oscillates between $-1/3$ and $1/3$. Note that the electrical conductance $G_\uparrow+G_\downarrow$ in the junctions is always positive and therefore spin current reversal in our model is not accompanied by current reversal.

The results in this limiting case are shown in Fig. 2. Figure 2(a) shows spin-resolved conductances as a function of χ, which is tunable by the gate voltage. Here, the phases of $G_\uparrow$ and $G_\downarrow$ are shifted by a half period, $\chi_+ - \chi_- = \pi/2$. Thus we have a finite spin current as shown in Fig. 2(b). Remark-

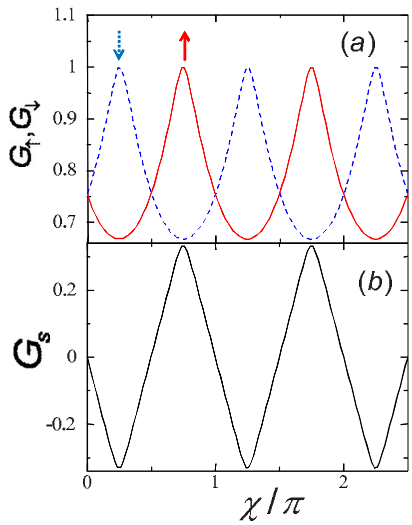

FIG. 2. (Color online) Conductance as a function of χ tunable by the gate voltage, in the limit of $U\to\infty$. (a) Spin-resolved conductances where the phases of $G_\uparrow$ (solid line) and $G_\downarrow$ (dotted line) are shifted by a half period, $\chi_\uparrow-\chi_\downarrow=\pi/2$. (b) Spin conductance G_s, which oscillates with the period π with respect to χ but is never damped.

ably, it oscillates with the period π with respect to χ but is never damped. As is seen, we can reverse the spin current by changing the gate voltage. Here, we focus on the limit of $|U\pm H|\gg E_F$. Next, we consider more general cases without taking this limit.

Figure 3(a) exhibits the spin conductance G_s as a function of χ for several values of k_FL. For a ferromagnetic graphene with $k_FL=0.1$, π periodicity is seen. With increasing k_FL, however, the π periodicity is broken as seen in Fig. 3(a). This is because $U\gg E_F$ is no longer satisfied for large k_FL since $\chi=k_FLU/E_F$. Thus, by choosing large χ for large k_FL,

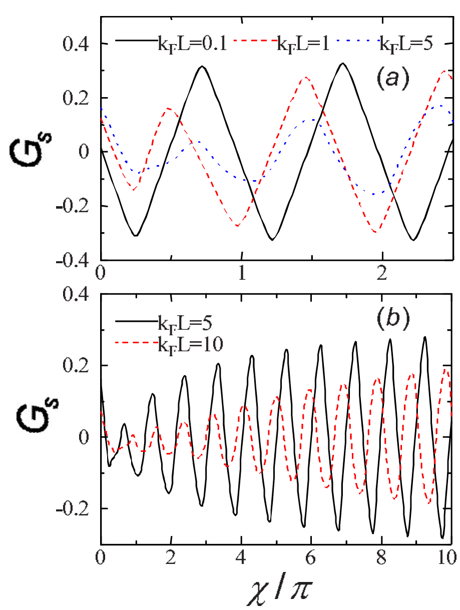

FIG. 3. (Color online) Spin conductance G_s as a function of χ for several values of k_FL with $\chi_+-\chi_-=\pi/2$.

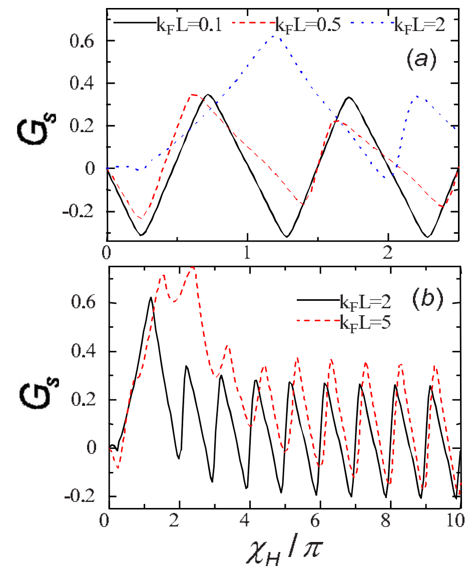

FIG. 4. (Color online) G_s as a function of χ_H, which is proportional to the exchange field, for several values of k_FL. Here, we choose $\chi=\pi/4$.

we again get π periodicity of G_s with respect to χ as shown in Fig. 3(b). From this figure, we find that $U/E_F>1$ is required to obtain controllable spin current reversal.

Figure 4 displays G_s as a function of χ_H, which is proportional to the exchange field, for several values of k_FL. Here, we fix $\chi=\pi/4$ to obtain a finite spin current. For a short ferromagnetic graphene with $k_FL=0.1$, π periodicity is satisfied. With increasing k_FL, the π periodicity requires a larger magnitude of χ_H as shown in this figure, similar to Fig. 3.

Our prediction is applicable as long as the continuum Dirac equation is valid, which requires a wide graphene nanoribbon. For the realization of our prediction of controllable spin current by the gate voltage, $U/E_F>1$ is required. If we choose ferromagnetic graphene with $k_FL=1$ and $E_F\simeq 1$ meV as an example, we need a length around 1 μm. Also, $\chi\sim U/E_F$, $\chi_H\sim H/E_F$, and hence U, $H\sim 1-10$ meV is required. These values can be achieved by the present experimental technique.[7–9,18,19]

In summary, we studied spin transport in normal/ferromagnetic/normal graphene junctions where a gate electrode is attached to the ferromagnetic graphene. We found that, because of the exchange field in the ferromagnetic graphene, the spin current through the junctions has an oscillatory behavior with respect to the chemical potential in the ferromagnetic graphene, which can be tuned by the gate electrode.[7–9] In particular, spin current reversal by the gate voltage, which is not accompanied by current reversal, is obtained. The exchange field of graphene is also known to be tunable by the in-plane external electric field.[18,19] Therefore, our prediction of high controllability of spin transport in ferromagnetic graphene junctions will contribute to the development of spintronics.

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[1] T. Ando, J. Phys. Soc. Jpn. **74**, 777 (2005).
[2] M. I. Katsnelson and K. S. Novoselov, Solid State Commun. **143**, 3 (2007).
[3] A. H. Castro Neto, F. Guinea, N. M. R. Peres, K. S. Novoselov, and A. K. Geim, arXiv:0709.1163v1 (unpublished).
[4] G. E. Volovik, *The Universe in a Helium Droplet* (Clarendon, Oxford, 2003).
[5] A. M. Tsvelik, *Quantum Field Theory in Condensed Matter Physics* (Cambridge University Press, Cambridge, U.K., 2007).
[6] K. S. Novoselov, A. K. Geim, S. V. Morozov, D. Jiang, Y. Zhang, S. V. Dubonos, I. V. Grigorieva, and A. A. Firsov, Science **306**, 666 (2004).
[7] Y. Zhang, Y.-W. Tan, H. L. Stormer, and P. Kim, Nature (London) **438**, 201 (2005).
[8] K. S. Novoselov, A. K. Geim, S. V. Morozov, D. Jiang, M. I. Katsnelson, I. V. Grigorieva, S. V. Dubonos, and A. A. Firsov, Nature (London) **438**, 197 (2005).
[9] K. S. Novoselov, E. McCann, S. V. Morozov, V. I. Fal'ko, M. I. Kastenelson, U. Zeitler, D. Jiang, F. Schedin, and A. K. Geim, Nat. Phys. **2**, 177 (2006).
[10] K. Yang, Solid State Commun. **143**, 27 (2007).
[11] H. B. Heersche, P. Jarillo-Herrero, J. B. Oostinga, L. M. K. Vandersypen, and A. F. Morpurgo, Nature (London) **446**, 56 (2007).
[12] C. L. Kane and E. J. Mele, Phys. Rev. Lett. **95**, 226801 (2005).
[13] D. Huertas-Hernando, F. Guinea, and A. Brataas, Phys. Rev. B **74**, 155426 (2006).
[14] N. Tombros, C. Jozsa, M. Popinciuc, H. T. Jonkman, and B. J. van Wees, Nature (London) **448**, 571 (2007).
[15] M. Ohishi, M. Shiraishi, R. Nouchi, T. Nozaki, T. Shinjo, and Y. Suzuki, Jpn. J. Appl. Phys., Part 2 **46**, L605 (2007).
[16] S. Cho, Yung-Fu Chen, and M. S. Fuhrer, Appl. Phys. Lett. **91**, 123105 (2007).
[17] H. Haugen, Daniel Huertas-Hernando, and Arne Brataas, arXiv:0707.3976v1 (unpublished).
[18] Y.-W. Son, M. Cohen, and S. G. Louie, Nature (London) **444**, 347 (2006).
[19] Er-Jun Kan, Zhenyu Li, Jinlong Yang, and J. G. Hou, arXiv:0708.1213v1 (unpublished).
[20] S. Dutta and S. K. Pati, arXiv:0712.4130v1 (unpublished).
[21] J. González, F. Guinea, and M. A. H. Vozmediano, Phys. Rev. B **59**, R2474 (1999).
[22] C. L. Kane and E. J. Mele, Phys. Rev. Lett. **93**, 197402 (2004).
[23] A. F. Morpurgo and F. Guinea, Phys. Rev. Lett. **97**, 196804 (2006).
[24] P. R. Wallace, Phys. Rev. **71**, 622 (1947).
[25] M. I. Katsnelson, K. S. Novoselov, and A. K. Geim, Nat. Phys. **2**, 620 (2006).
[26] S. Bhattacharjee and K. Sengupta, Phys. Rev. Lett. **97**, 217001 (2006).
[27] J. Linder and A. Sudbø, Phys. Rev. Lett. **99**, 147001 (2007).