Spin polarization and magnetoresistance through a ferromagnetic barrier in bilayer graphene

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Abstract

We study spin dependent transport through a magnetic bilayer graphene nanojunction configured as a two-dimensional normal/ferromagnetic/normal structure where the gate voltage is applied on the layers of ferromagnetic graphene. Based on the four-band Hamiltonian, conductance is calculated by using the Landauer–Buttiker formula at zero temperature. For a parallel configuration of the ferromagnetic layers of bilayer graphene, the energy band structure is metallic and spin polarization reaches its maximum value close to the resonant states, while for an antiparallel configuration the nanojunction behaves as a semiconductor and there is no spin filtering. As a result, a huge magnetoresistance is achievable by altering the configurations of ferromagnetic graphene around the band gap.

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

1. Introduction

Since spin–orbit coupling in graphene [1] is very weak [2] and there is no nuclear spin [3], the spin flip length is very long, about 1 µm, in dirty samples and at room temperature [4]. Clean samples are expected to have a longer spin coherency, and this is a good opportunity develop spintronic applications based on graphene. On the other hand, graphene has no intrinsically ferromagnetic (FM) properties, but it is possible to induce ferromagnetism externally by doping and defects [5], Coulomb interactions [6] or by applying an external electric field in the transverse direction in nanoribbons [7]. Recently, Haugen [8] proposed FM correlations as being due to the presence of magnetic states close to graphene. The overlap between the wavefunctions of the localized magnetic states in the magnetic insulator and the itinerant electrons in graphene induces an exchange field on itinerant electrons in graphene, giving rise to spin splitting of the transport. The exchange splitting induced by the FM insulator EuO in graphene was estimated to be of the order of 5 meV. This splitting, which is effectively similar to a Zeeman interaction, has such a large magnitude that it can have important effects. Such spin splitting can be directly evaluated from the transmission resonant peaks or magnetoresistance of the FM graphene junction. The ferromagnetism leads to spin splitting effectively similar to a Zeeman interaction but of a much larger magnitude. The induced exchange field is tunable by an in-plane external electric field [9]. The possibility of controlling spin conductance in a FM graphene monolayer insulator has also been studied by Yokoyama [10]. It was found that the spin conductance has an oscillatory behavior in terms of chemical potential and the gate voltage.

Bilayer graphene, on the other hand, has been shown to have interesting properties for application in nanoelectronic devices such as transistors based on a graphene substrate. A new type of integer quantum Hall effect [11] and also an electronic band gap controllable by a vertically applied electric field are among its unusual properties compared to monolayer graphene [12–15]. Moreover, the parabolic band structure close to the Dirac points transforms to a sombrero-like dispersion when an electric field is applied to the graphene. Optical measurements and theoretical predictions propose a 200 meV gap in bilayer graphene. This controllable gap makes bilayer graphene an appropriate candidate for spintronic devices. An effective two-band Hamiltonian can describe the low energy excitations of a
graphene bilayer in the regime of low barrier heights [16]. However, a four-band Hamiltonian is known to give a better agreement with both experimental data and theoretical tight-binding calculations [12, 17]. Very recently, spin splitting of conductance in bilayer graphene was investigated by using an effective two-band Hamiltonian emerging from a low energy approximation [18]. This approximation is valid when the energy of electrons hitting the potential barrier approximates the barrier height. Application of a potential difference between the upper and lower layers intensifies the failure of this approximation. On the other hand, magnetoresistance in bilayer graphene has been studied in [19] by using an 8 × 8 Hamiltonian when the induced exchange fields are laid in the plane of each layer with a rotation in their orientations against each other. By using the Kobo–Greenwood formula, they investigated the dependence of conductivity and magnetoresistance on temperature and induced exchange field.

Motivated by these studies, based on the four-band Hamiltonian and close to the Dirac points, we study spin current through a magnetic barrier created by the proximity of a ferromagnetic insulator on bilayer graphene. Conductance is calculated by use of the Landauer–Buttiker formula at zero temperature. The parameters of the barrier, energy and angle of the incident electrons can affect transport through a magnetic barrier, classified into propagating or evanescent modes. The dependence of resonant peaks in transmission on the system parameters is proposed to follow a resonance condition. We have found in some resonant energies and barrier parameters, and also around the gaped region, that a remarkable spin polarization and also magnetoresistance can be achieved.

For our model we consider a normal/ferromagnetic/normal bilayer graphene nanojunction. The model we have used for bilayer graphene sandwiched between two ferromagnetic insulators is shown schematically in figure 1. The exchange fields induced by the ferromagnetic insulators is supposed to be perpendicular to the graphene plane. Therefore, the Hamiltonian of the spin up detaches from the spin down. Two gate electrodes can be attached to the ferromagnetic graphene from the upper and lower layers which control the barrier height in each layer. This setup is different from the systems studied by [19, 20]. The exchange field splits this potential depending on the orientation of the exchange field. So in the ferromagnetic part of bilayer graphene, we have \( V_0 = V_0 \pm \Delta \), where \( \Delta \) and \( V_0 \) are the exchange field and the potential barrier made by the gate voltage, respectively. So the two spins are scattered from the barriers with different heights. This means that energy shift of the top of the valance band in the barrier is different for parallel and antiparallel spins to the exchange field. This spin splitting causes a conductance shift as a function of energy for each spin, resulting in magnetoresistance. To investigate spin polarization and also magnetoresistance, we have considered two different configurations so that the exchange field inducing by the magnetic insulators on each layer are parallel or antiparallel with respect to each other. The configuration prepared for observation of magnetoresistance differs from the configuration considered by [19, 20]. The parallel configuration has a metallic behavior, while the antiparallel configuration induces a potential difference between upper and lower layers; we conclude that the system has a semiconductor behavior with a band gap of \( 2\Delta \).

This paper is organized as follows: we briefly explain the formalism which is used for calculating the transmission based on the four-band Hamiltonian. Before we present our results, it is important to give a short review in section 3 on transport through a barrier deposited on bilayer graphene and its dependence on the system parameters such as energy of incident quasi-particles and their angle of incidence on the barrier and also barrier parameters. The method presented in section 2 is a detailed analysis accompanied by some small corrections to the method used by [17]. We will present spin polarization in the parallel configuration in section 4. Magnetoresistance and its dependence on the energy of incident particles and also the induced magnetic field will be investigated in section 4. Finally, we give our conclusion in section 5.

2. Formalism

In the unit cell of bilayer graphene, we suppose that two independent sublattices, A and B, related to each monolayer of graphene are connected to each other in the Bernal stacking. Close to the Dirac points and in the nearest neighbor tight-binding approximation, the four-band Hamiltonian and also its eigenfunction is written as the following:

\[
H = \begin{pmatrix}
V_1 & \pi & t_\perp & 0 \\
\pi^\dagger & V_1 & 0 & 0 \\
t_\perp & 0 & V_2 & \pi^\dagger \\
0 & 0 & \pi & V_2 \\
\end{pmatrix}, \quad \Psi = \begin{pmatrix}
\psi_A \\
\psi_B \\
\psi_F \\
\psi_A' \\
\end{pmatrix}
\]

(1)

where

\[
\pi = (p_x + ip_y)v_F = -i\hbar v_F(\partial_x - k_y)
\]

and in the above formula \( k_y = k \sin \theta \) where \( \theta \) and \( k \) are the incident angle and wavenumber of quasi-particles hitting a barrier which is created by applying a bias to a metallic strip deposited on bilayer graphene. Moreover, \( V_1 \) and \( V_2 \) are the gate potentials applied to the upper and lower layers of bilayer graphene. Such a gate potential can be manipulated by applying a perpendicular electric field on a graphene sheet.

Figure 1. Schematic view of normal/ferromagnetic/normal bilayer graphene. Two gate electrodes can be coated on top of the magnetic insulator strips which are located on the upper and lower layers.
Here, the barrier is approximated by a square potential barrier with sharp variation. By solving the eigenvalue equation of $H\Psi = E\Psi$, the four-band spectrum can be concluded to be the following:

$$
\left(\epsilon'\right)^2 = k^2 + \delta^2 + \frac{\left(t'\right)^2}{2} \pm \sqrt{4k^2\delta^2 + \left(t'\right)^2 + \left(\frac{t'}{2}\right)^2 + k^2}
$$

where the above parameters are defined as the following:

$$
\epsilon' = \left(E - V_0\right)/\hbar v_F = \epsilon - \nu_0, \quad V_0 = \left(V_1 + V_2\right)/2
\delta = \left(V_1 - V_2\right)/2\hbar v_F, \quad t' = t/\hbar v_F.
$$

In the case of $\delta = 0$ and $k \ll t'$, in the low energy limit, the energy spectrum behaves as $E - V_0 = \pm \hbar^2 k^2/2m$, where $m = t/(2\hbar v_F)$ is an effective mass. This approximation which results in an effective two-band Hamiltonian is valid when the energy of incident electrons is close to the barrier height. In other words, in the case of zero potential difference between two layers, the absolute value of $(E - V_0)$ should be much smaller than the interlayer coupling strength (0.4 eV) to preserve smallness of the wavenumber inside the barrier. Moreover, to have a small wavenumber outside the barrier, the energy of incident electrons has to be small. So in this approximation barrier height and energy also have to be small. For $\delta \neq 0$, this approximation may fail for large potential differences. So one would wish to choose valid energy and potential ranges. However, in spite of [18], in this paper we use a four-band Hamiltonian in which the only approximation is the Dirac cone.

If we assume a plane wave solution for the Schrödinger equation, the wavefunction in each region with a constant potential is written as the following matrix product:

$$
\Psi = GM\begin{pmatrix} a \\ b \\ c \\ d \end{pmatrix}
$$

where the matrix elements of the matrices $G$ and $M$ are defined as:

$$
G = \begin{pmatrix}
1 & 1 & 1 & 1 \\
-f_+ & f_+ & f_+ & f_+ \\
h^+ & h^+ & h^+ & h^+ \\
g_+^h & g_+^h & g_+^h & g_+^h \\
\end{pmatrix}
$$

$$
M(x) = \begin{pmatrix}
e^{i\alpha(\pm x)} & 0 & 0 & 0 \\
0 & e^{-i\alpha(\pm x)} & 0 & 0 \\
0 & 0 & e^{i\alpha(\pm x)} & 0 \\
0 & 0 & 0 & e^{-i\alpha(\pm x)} \\
\end{pmatrix}
$$

$$
\begin{aligned}
f_\pm &= \frac{\pm \alpha_+ - ik_x}{\epsilon' - \delta} & f_\pm &= \frac{\pm \alpha_- - ik_x}{\epsilon' - \delta} \\
g_\pm &= \frac{\pm \alpha_+ + ik_x}{\epsilon' + \delta} & g_\pm &= \frac{\pm \alpha_- + ik_x}{\epsilon' + \delta} \\
h_\pm &= \frac{(\epsilon' - \delta)^2 - \alpha_+^2 - \alpha_-^2}{t'(\epsilon' - \delta)} \\
\end{aligned}
$$

Here $\alpha$ is the wavenumber in the $x$ direction and is defined as:

$$
\alpha^2 = \left[\delta^2 + \left(\epsilon'\right)^2 \right] - k_x^2 \pm \sqrt{4\left(\epsilon'\right)^2 + \left(t'\right)^2 + \left(\frac{t'}{2}\right)^2 + k_x^2}.
$$

In the special case of normal incident angle and zero gate potential, where $k_x = \delta = 0$, the wavenumber $\alpha_+$ is real in the energy range of $0 < \epsilon' < t'$ and $\alpha_-$ is real if the energy of incident particles is in the range of $\epsilon' < 0$, $\epsilon' > -t'$. In this paper, the studied system contains a magnetic or electrostatic potential barrier as shown in figure 2. The barrier width is $w$. The electrostatic potential which plays the role of the gate voltage is set to be $V_0$ in the barrier part and zero in the first and last regions. We suppose that the energy range of incident particles is limited to the range of $0 < \epsilon' < t'$. Consequently in the barrier part we have $-V_0 < \epsilon' < t' - V_0$.

The wavenumbers behind and in front of the barrier $\alpha(1)_+$ and $\alpha(3)_+$ are real while $\alpha(1)_-$ and $\alpha(3)_-$ are imaginary. In the barrier part, for $\epsilon' < 0$, the wavenumbers $\alpha(2)_+$ and $\alpha(2)_-$ are imaginary and real, respectively, and vice versa for $\epsilon' > 0$. A schematic view of the barrier at normal incidence and wavenumbers in each part is shown in figure 2.

By applying continuity of the wavefunctions on the boundaries of the barrier, one can connect the coefficient matrix of the wavefunction for the last region $A_3$ to the coefficient matrix for the first region $A_1$:

$$
A_1 = NA_3
$$

$$
N = M_1^{-1}(0)G_1^{-1}G_2M_2^{-1}(w)G_3^{-1}G_4M_3(w)
$$

where $N$ is called a transfer matrix. Since $\alpha(1)_-$ and $\alpha(3)_-$ are imaginary in the energy range of interest, those parts of the wavefunction that are associated with such wavenumbers are an exponentially growing or decaying function. So we have to set the coefficient of the plane wave $e^{i\alpha(1)_-x}$ (c in equation (4)) to be zero for the first region, because this part of the wavefunction grows exponentially when $x \to -\infty$. Therefore, the coefficient matrix in the first region is supposed to be $A_1 = [1, r, 0, e\bar{g}]^T$, where the superscript $T$ refers to the transpose of a matrix, $e\bar{g}$ is the coefficient of the growing evanescent state and $r$ is the coefficient of the reflected part of the wavefunction. In the last region, we have to set the coefficient (d in equation (4)) of $e^{-i\alpha(3)_+x}$ to be zero because this part of the wavefunction increases exponentially when
For energies $\epsilon > V_0$, transmission becomes complete ($T \rightarrow 1$) and so transmission is complete ($T \rightarrow 1$). By applying a vertical electric field in the barrier part, a band gap is opened in the band structure of bilayer graphene which is proportional to the potential difference between the potentials of each layer. In this case, chiral symmetry fails and therefore transmission in normal incidence is nonzero for energies lower than the barrier height ($E < V_0$). Transmission at normal incidence is represented in figure 3 as a function of $\epsilon'_{2}$ for $\delta = 0$ and $10$ meV. Application of a vertical electric field brings about the emergence of some resonant tunneling states for energies of $E < V_0$. In this energy range, resonant states originate from interference of the incident and scattered waves. For all cases such as nonzero incident angle and $\delta \neq 0$, the resonant peaks are interpreted by the proposed resonance condition relation:

$$\alpha_0 w = \pi \delta \sin \epsilon'_{2}.$$  

(13)

where $\alpha_0$ is the $x$-component of the wavenumber inside the barrier which can be calculated by equation (7).

To have a more complete view, we prepare a contour plot of transmission in the plane of the incident angle and $\epsilon'_{2}$, which is shown in figure 4 for a fixed barrier width. For normal incidence ($\theta = 0$), transmission behavior is compatible with the results shown in figure 3.

For energies higher than the barrier height $\epsilon'_{2} > 0$, transmitting channels are opened over the whole energy range. However, the transmitting window for the incident angles is limited by the condition that $\alpha'^{2}_{2}$ (in equation (7)) is real. In the case of $\delta = 0$, the range of incident angles in which transmission is high can be extracted as $\tan^{-1} \sqrt{\frac{\epsilon'_{2}^2 + \epsilon'_{2}'}{k}} \leq \theta \leq \tan^{-1} \sqrt{\frac{\epsilon'_{2}^2 + \epsilon'_{2}'}{k}}$. Therefore, by increasing $\epsilon'_{2}$, the range of angles with high transmission becomes more extended. In the energy range $\epsilon'_{2} < 0$, independent of the value of $\delta$, transmission is complete ($T \rightarrow 1$) and so transmission is complete ($T \rightarrow 1$). By applying a vertical electric field in the barrier part, a band gap is opened in the band structure of bilayer graphene which is proportional to the potential difference between the potentials of each layer. In this case, chiral symmetry fails and therefore transmission in normal incidence is nonzero for energies lower than the barrier height ($E < V_0$). Transmission at normal incidence is represented in figure 3 as a function of $\epsilon'_{2}$ for $\delta = 0$ and $10$ meV. Application of a vertical electric field brings about the emergence of some resonant tunneling states for energies of $E < V_0$. In this energy range, resonant states originate from interference of the incident and scattered waves. For all cases such as nonzero incident angle and $\delta \neq 0$, the resonant peaks are interpreted by the proposed resonance condition relation:

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resonant peaks emerge for nonzero incident angles ($\theta \neq 0$) which obey the resonance condition $\alpha_{0}w = n\pi$. So in addition to some resonant energy states, we have some resonant widths in which transmission is high. Figure 5 shows transmission in the plane of the incident angle and the barrier width for $\varepsilon'_{2} < 0$ and $\delta = 0$. It is shown that, based on the resonance condition (equations (13) and (7)), for large incident angles $\alpha_{0}$ is reduced and so in a fixed resonance order ($n$) the resonance condition is satisfied for wide barriers. Therefore, the resonance strips with complete transmission shown in figure 5 depend strongly on the incident angle in the range of wide barriers.

By applying a vertical electric field in the barrier part, a band gap is opened around $\varepsilon'_{2} = 0$. This band gap also has a trace in transmission as a transport gap, shown in figure 4(b).

4. Results

By application of an averaged gate voltage $V_{0}$, band structure in the barrier part is shifted by the value of $V_{0}$. Figure 6 shows the band structure of magnetic insulators with parallel and antiparallel configurations when a gate voltage is applied on the barrier part. If the exchange fields induced in each layer of bilayer graphene are parallel, particles with spin parallel (spin up) and antiparallel (spin down) to the exchange fields are scattered from barriers with different heights. In the parallel configuration, spin splitting of the barrier potential in the ferromagnetic graphene is written as $V^{-} - V^{+} = 2\Delta$. Such spin splitting is also seen in the band structure that is shown in figure 6(a). It is seen that the top of the valence band is shifted to lower/higher energies for spins up/down. However, in the antiparallel configuration, the band structure shown in figure 6(b) is the same for both up and down spins. A band gap which is proportional to $2\Delta$ appears in the band structure of the antiparallel configuration.
4.1. Spin polarization

Here, there is a correspondence between the band structure and transmission. According to the band structure, we expect to have spin polarization to emerge just for the parallel configuration because energy bands for up and down spins are shifted by $2\Delta$ with respect to each other. However, since the band structure for the antiparallel configuration is the same for both spins, it is not expected to have spin polarization for this configuration. The spin polarization is defined as:

$$P = \frac{G_{\text{up}} - G_{\text{down}}}{G_{\text{up}} + G_{\text{down}}}$$

where $G_{\text{up}}$ and $G_{\text{down}}$ are the conductance for up and down spins, respectively. The conductance is calculated by using Landauer formalism in the linear regime. Therefore, conductance is proportional to angularly averaged transmission projected along the current direction.

$$G = \int_{-\pi/2}^{\pi/2} T(E, \cos(\theta)) \cos \theta \, d\theta.$$  

It is clear that additional to the transmission curves (figure 4), resonance peaks also appear in conductance. Since up and down spins in the parallel configuration see barriers of different heights, resonance peaks in conductance as a function of Fermi energy $E$ are shifted to higher energies as $\Delta$ for spin down and to lower energies as $-\Delta$ for spin up. This mismatching of conductance peaks for two spins causes a large spin polarization at resonance states. Figure 7 displays conductance and spin polarization as a function of $\varepsilon'_2$ for the parallel configuration. It is shown that conductance peaks, and consequently spin polarization, appear in the energy range of $\varepsilon'_2 < 0$. It is seen that by inducing an exchange field, conductance peaks in figure 7(a) split into two peaks which are related to each spin. This spin splitting is about $2\Delta$. Spin polarization shown in figure 7(b) has an oscillatory behavior with energy of incident particles for energies lower than the barrier height $\varepsilon'_2 < 0$. The amplitude of spin polarization increases with the induced exchange field $\Delta$ and reaches its maximum value. However, spin polarization tends to zero for energies greater than the potential height $\varepsilon'_2 > 0$ except at $E \sim V_0$.

In the parallel configuration and for $\varepsilon'_2 < 0$, figure 8(a) shows that conductance in the resonance widths has a peak. These peaks, which are also seen in the transmission curves of figure 5, are explained by the resonance condition of equation (13). It is shown that spin splitting of conductance peaks also appears in the resonance widths and originates from the different barrier heights for two spins up and down. It should be noted that the conductance at resonance widths decreases for wide barriers. In the wide range of widths, the angular window for transmitting channels shown in figure 5 decreases with the width.

Figure 8(b) shows spin polarization as a function of the barrier width. Again, spin polarization has an oscillatory behavior with barrier width. The amplitude of spin
polarization strongly increases with an increase in the induced exchange field. Therefore, to manifest this spin polarization, we should manufacture the ferromagnetic graphene part in the particular widths in which spin polarization reaches the value of unity.

4.2. Magnetoresistance

In this section we will show that by switching between parallel and antiparallel configurations one can obtain large magnetoresistance. Magnetoresistance is defined as follows:

\[
\text{MR} = \frac{G^p - G^{ap}}{G^p + G^{ap}}
\]

where \(G^p = G^p_{\text{up}} + G^p_{\text{down}}\) and \(G^{ap} = G^{ap}_{\text{up}} + G^{ap}_{\text{down}}\) are conductance for parallel and antiparallel configurations.

Figure 9 displays conductance in the parallel and antiparallel configurations and also magnetoresistance as a function of \(\varepsilon'_2 = (E - V_0)/\hbar v_F\) for a barrier with a width of 40 nm and the barrier width for a barrier with a height of 50 meV. As we stated before, spin splitting at the resonance states (for \(\varepsilon'_2 < 0\)) emerges in conductance peaks in the case of the parallel configuration. This behavior is clear in figures 9(a) and (c). However, this splitting will not occur for the case of an antiparallel configuration. Therefore, large magnetoresistance appears around the conductance resonance peaks. In the parallel configuration, a band gap appears around the barrier edge in the interval \(V_0 - \Delta < E < V_0 + \Delta\).
This band gap has a trace in transmission, and consequently conductance. The zero conductance region around the barrier edge $\varepsilon_2^0 \sim 0$, which is seen in figure 9(a), is a result of the band gap. Since there is no such band gap in the parallel configuration, magnetoresistance as shown in figure 9(b) reaches its maximum value in the energy band gap. In an energy range greater than the barrier height $\varepsilon_2^0 > 0$, there is no spin splitting, and therefore magnetoresistance tends to zero.

As we showed before, conductance has peak at resonant widths. Similar to the previous case, spin splitting occurs just for the parallel configuration. So magnetoresistance increases around the resonance widths. The oscillatory behavior of magnetoresistance as a function of the barrier width is represented in figure 9(d).

As we showed, there is a large magnetoresistance around the barrier edge $E \approx V_0$. In this range of energies, we investigate the dependence of magnetoresistance on the induced exchange field. This exchange field of graphene can be controlled by an in-plane external electric field [9]. Figure 10(b) represents that magnetoresistance increases monotonically by increasing the exchange field $\Delta$. It is interesting that by increasing the exchange to 10 meV, magnetoresistance reaches its maximum value.

To explain this behavior, we investigate the dependence of conductance on the exchange field in the parallel and antiparallel configurations. In the antiparallel configuration, the band gap, which is limited in the interval of $V_0 - \Delta < E < V_0 + \Delta$, is enhances by increasing the exchange field. Therefore, conductance in the antiparallel configuration goes to zero when the exchange field is increased. Suppression of the conductance with the exchange field in the antiparallel configuration is shown in figure 10(a). However, in the parallel configuration, conductance increases by enhancement of the exchange field. The reason for this enhancement is the larger angularly transmitting windows for larger $\varepsilon_2^0$ (see figure 4). In fact, the effective potential for spins up $\hat{V}^+ = V_0 - \Delta$ is decreased by an increase in $\Delta$. So $\varepsilon_2^0 = (E - V_0)/\hbar v_F$ for a fixed energy is increased and consequently $G^{up}$ and so $G^p$ is increased by $\Delta$. As a conclusion, for exchange fields up to 10 meV, suppression of $G^p$ and an increase in $G^p$ results in a large magnetoresistance which is very useful for designing spin memory devices.

5. Conclusion

We have studied spin polarization and magnetoresistance of a normal/ferromagnetic/normal junction of bilayer graphene by using a transfer matrix method based on the four-band Hamiltonian. Transport properties are simultaneously controlled by two gate electrodes ($V_0$), which are applied on the ferromagnetic graphene. Two configurations of the exchange field are considered perpendicular to the graphene sheet. This exchange field is induced by the proximity of a localized magnetic orbital in a magnetic insulator coating on top of each layer of bilayer graphene in the barrier part. In the parallel configuration in which graphene has a metallic behavior, a spin splitting $2\Delta$ occurs for the conductance at the resonant states just for energies lower than the barrier height $E < V_0$. However, there is no spin splitting in the antiparallel configuration. A band gap of $2\Delta$ is opened in the antiparallel configuration which makes it a semiconductor. As a result of spin splitting in the parallel configuration, an oscillating spin polarization emerges for energies lower than the barrier height. Furthermore, an oscillatory magnetoresistance with large amplitude is achievable for $E < V_0$ when we are able to switch between two configurations. There is also a large magnetoresistance in the energy range around the barrier edge. In this range of energy, magnetoresistance reaches its maximum value when the exchange field is increased by an in-plane external electric field.

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