Magnetoresistance effect realized in current-in-plane Van der Waals spin valve structure by electrically switchable magnetization

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Abstract
We present a device concept for magnetoresistance (MR) effect in which two Van der Waals (vdW) spin valves are used as the left and right ferromagnetic (FM) leads connected by a bilayer graphene as the channel material of the central scattering region. Unlike conventional current-perpendicular-to-plane magnetic tunnel junction consisting of two FM thin films with a thin insulating barrier sandwiched between them, the FM leads in our proposed current-in-plane MR device are vDW spin valves. This is important because the application of an out-of-plane electric field allows control of the direction and magnitude of the magnetization in vDW spin valves. Moreover, we show that the oscillatory behavior is found in the MR and conductance as the height (depth) of the barrier (well) of the scattering region with the p(n)-doping increase, or the width of the scattering region increase. Remarkably, when the potential barrier is present, the oscillation magnitude of MR is considerable and can be changed from positive to negative value, whereas for the potential well, the oscillation amplitude is relatively small and is always positive. Therefore, we hope that this device configuration with electrically tunable large MR can open up new possibilities for future lower power magneto-electric devices since no current-induced magnetization switching.

1. Introduction
Since the discovery of magnetoresistance (MR) in a Fe/Cr layered system [1], the effect of MR enables the present information technology age. Research in conventional MR effect device, which has two ferromagnetic (FM) layers sandwiching a thin insulating barrier, has been widely used in magnetically memory storage and magnetosensing technologies [2–4]. From the MR device point of view, it is a crucial issue to find most efficient paths for switching two magnetic layers from parallel (P) to antiparallel (AP) alignments of the magnetization. Research in traditional approach for manipulating magnetization is to use electric current to generate a large magnetic field. Spin-transfer torque generated by spin polarized current is also a promising approach for manipulation of magnetization [5, 6]. Recent discoveries regarding current-induced spin–orbit torques produced by materials with strong spin–orbit interactions dramatically improved efficiency in the manipulation of magnetization [7–10].

Current-induced magnetization switching, however, inevitably involves Joule heating, which significantly increase the switching power. For low power and high-speed magnetic switching applications, the electric-field manipulation of magnetism has sparked intense research interest in condensed matter physics [11], such as thin-film ferromagnets [12, 13], FM transition-metal monolayers [14], multiferroics [15], magnetic tunnel-junctions [16], van der Waals (vdW) magnetic semiconductors [17, 18]. Although the idea that an electric field through a voltage could induce the reversal of ferromagnets has been around
since the 1960s [19], many studies have been focused only on the realization of electric-field-assisted reversible switching in the intrinsic magnetic materials including magnetic semiconductors, magnetic metals and multiferroic materials.

Another technological challenge, raised by the conventional current-perpendicular-to-plane MR effect device, is that the disordered interface between the FM layer/insulating barrier interface, which determines the efficiency of the conventional MR [20–22]. Recently, a suitable approach of obtaining high quality interface is replacing thin insulating barrier with atomically thin spacers, such as graphene and other two dimensional (2D) materials [23–26]. For example, experimentally, current-perpendicular-to-plane magnetic tunnel junction using 2D materials, such as graphene [27–29], MoS2 [30], and boron nitride [31] as a nonmagnetic spacer, have been fabricated successfully.

Note that this was attributed to the fact that 2D materials exhibit novel crystal structures and physical properties, which is different from their bulk counterpart [32–34]. For instance, it has been well known that monolayer graphene due to the chirality of the Dirac electron exhibits Klein tunneling, which predicts that the Dirac electron could completely pass through the step regardless of the step height at normal incidence, while massive Dirac electron in Bernal-stacked BLG shows perfect reflection injecting normal to the potential step [35]. For oblique incidence, the finite size of the central region results in the appearance of the characteristic Fabry–Pérot (FP) resonances in the transmission, which is a demonstration of the quantum interference between Dirac electron waves in monolayer graphene and BLG [35–37]. Experimentally, graphene-based FP interferometry have been used to observe standing waves at graphene edge [38], electronic thickness of graphene [39], fractional quantum Hall effect [40]. Moreover, the high carrier mobility, quantum Hall effect and ballistic transport of Dirac electron in 2D materials have opened the door to the Dirac-electron-based devices and phenomena, ranging from Mach–Zehnder oscillations [41] to Veselago lens [42, 43], and even the interference of d-wave superconducting pairs in graphene [44].

However, experimentally reported MR in current-perpendicular-to-plane MR device with 2D materials as spacer layers is also quite low, which is undesirable [27–31]. For example, the transport measurements in NiFe-graphene-Fe junction by Cobas et al [28] give a negative MR of −12% at 15 K and −5% at room temperature. Similar studies [27, 29, 30] also find that the highest tunnel MR up to 6% is observed, which do not show very high MR predicted for an ideal system. The low MR was attributed to the use of Permalloy electrodes, such as Fe, Ni, and Co, in which only minority spin-polarized electrons can transmit from the FM electrodes into 2D materials. Thus, it is important to search for different 2D materials stacks to inject current with a relatively high spin polarization. Specially, stacking 2D magnetic materials can offer better opportunities to use 2D material as FM electrodes. This is because that proximity with ferromagnets can induce changes in the electronic structure of such 2D materials. Due to magnetic proximity effect, graphene and graphene-like 2D materials can become magnetic, opening opportunities for acquiring the properties of its neighbors, which are not present naturally within a single 2D material.

One prominent example of such a magnetic proximity effect is exhibited in a vdW spin valve structure, where bilayer graphene (BLG) is sandwiched between two insulating ferromagnets (IFMs) [45]. In this work, Cardoso et al predicted that BLG experiences strong spin splitting and band gap when the relative magnetization orientations of the two IFMs are aligned P and AP, respectively. Furthermore, it should be mentioned that BLG is also highly desirable for device applications because of the electric-field induced band gap engineering [46], the more precise control of the chemical potential [47], as well as the long spin diffusion lengths and spin lifetimes [48, 49]. Therefore, there is special interest in ways to modulate the proximity-induced magnetic properties in BLG-based vdW spin valves by means of an electric field, allowing for an ultraefficient electric-field control of magnetoresistive random-access memories [50]. To date, several methods have been used to realize electric field control of magnetic properties [12, 13, 51–53], but further breakthroughs are required.

In this study, we demonstrate the existence of current-in-plane MR effect in a BLG structure where vdW spin valves are used as the FM leads. This is important because we can manipulate the direction and magnitude of the magnetization in the vdW spin valves by out-of-plane electric field. It is demonstrated that the MR and conductance in such a system is oscillatory with increasing the height (depth) of the barrier (well) of the scattering region in case of the p(n) doping, or the width of the scattering region. When the potential barriers (n − p − n) or wells (n − n′ − n, the prime referring to the scattering region) is constructed by controlling the doping in the scattering region, we find that the oscillation magnitude of MR is determined by the height (depth) of the barrier (well). Namely, with variations of the potential barrier, not only the magnitude of the MR can change, but also its sign can switch from positive to negative. For the potential well, however, a significant change of the MR is that the oscillation amplitude is always positive. Our findings show that this device configuration with electrically tunable large MR can open up new possibilities for lower power magneto-electric devices since no current-induced magnetization switching.
Figure 1. Device structure and band structures. (a)–(c) Schematic of the IFM encapsulated BLG device. The IFM/BLG/IFM vdW structure rests on a SiO2/Si substrate, which serves as a bottom gate, and Au is used as the top gate electrode, with the top $h$-BN layer working as the gate dielectric layer. Calculated band structure (in continuum approximation, near the momentum valleys) for the electric field: (c) negative $E_z \downarrow$ and (d) positive $E_z \uparrow$. Black and blue (red and magenta) colors represent spin-up and spin-down branches of the conduction(valance) band, respectively.

The remainder of the paper is laid out as follows. In section 2 we describe the theory and model we employ to calculate the MR effect of the BLG-based spin valve structure. In section 3, we show the numerical results and present our discussions. Our concluding remarks follow in section 4.

2. Model and theory

To study the electric field effect on the modulation of proximity-induced magnetic property in BLG-based vdW spin valves, the dual-gating scheme is under consideration [54–57], as shown in figures 1(a) and (c). In this scheme, a vdW stacking technique is used to encapsulate a AB-stacked BLG sheet by two IFMs, such as EuO [58, 59] and CrI3 [45]. This stack is then placed on a SiO2/Si substrate serving as the bottom gate. On top we place one piece of Au, which was used as the top gate electrode with the hexagonal boron nitride ($h$-BN) working as dielectric layer. This dual-gating scheme generate two effects: the average of the two voltages induces total carrier doping, that is, the Fermi energy and their difference leads to interlayer electric field $E_z$, that is, the band gap. With this arrangement, the low-energy effective Hamiltonian for BLG can be written as

$$H = H_0 + \lambda_E \sigma_z + s \lambda_{AFM} \sigma_z,$$

(1)

where the first term $H_0$ is the kinetic energy. The second term $\lambda_E = 2\ell E_z$ with perpendicular electric field $E_z$ and the interlayer distance $\ell$ is known to open up a gap, and moderate gate voltages allow the tuning of the band gap. Notice that in contrast to the conventional parallel plate capacitor, the difference is the case where the inserted metal plates are replaced with a 2D BLG. The classical electrostatics will not be valid anymore. Here an additional capacitive effect defined as a ‘quantum capacitance effect’ is found, which is a property of the many-body system. And this effect requires the associated electric field can penetrate through the 2D BLG [60–62]. Hence, the interlayer electric field $E_z$ in the BLG, as shown in figures 1(a) and (c), is not zero but finite. The third term $\lambda_{AFM} = \Delta$ with $\Delta = 3.5$ meV can also induce a band gap in the AP alignment of magnetizations in the IFM/BLG/IFM vdW spin valve structure [45]. $\sigma_z$ is the Pauli matrix. Further, $s = \pm 1$ represents spin-up ($\uparrow$) and spin-down ($\downarrow$) states.
Figure 2. Schematic view of our device and band structures. (a) The device structure composed of a scattering region and two semi-infinite electrodes. The left electrode and right electrode are vdW spin valves. (b) and (c) Schematic band structures of the left electrode, scattering region, and right electrode in the negative $E_z \downarrow$ (left) and positive $E_z \uparrow$ (right) electric field configurations. The horizontal dashed line denotes the Fermi energy $E_F$. $U$ is the value of the electrostatic potential at the scattering region. If the Fermi energy $E_F$ of the Dirac electrons is smaller (larger) than the electrostatic potential $U$, an $n−p−n(n′−n)$ junction forms.

The Hamiltonian $H_0$ in equation (1) is the effective two-band Hamiltonian of BLG [42, 63, 64], which is given by

$$H_0 = \begin{pmatrix} 0 & \frac{\hbar^2}{2m} (k_y + i\eta k_x)^2 \\ \frac{\hbar^2}{2m} (k_y - i\eta k_x)^2 & 0 \end{pmatrix},$$

(2)

with the effective mass $m = \frac{\gamma_1}{2v_F^2}$. Here, $\gamma_1 \approx 0.4$ eV is the interlayer coupling, $v_F \approx 10^6$ m s$^{-1}$ is the Fermi velocity. $\vec{k} = (k_x, k_y)$ is the 2D wave vector in the vicinity of the two momentum valleys. $\eta = \pm$ denotes two independent momentum valleys.

The energy dispersion associated with the Hamiltonian in equation (1) is given by

$$E_{\varepsilon,s} = \sqrt{\left(\frac{\hbar^2 k^2}{2m}\right)^2 + (\lambda_E + s\Delta)^2},$$

(3)

Here $\varepsilon = \pm$ specifies the conduction and valance bands. $k = \sqrt{k_x^2 + k_y^2}$ is the wave vector measured from the momentum valleys $K (K' \prime)$ point. From figures 1(b) and (d), we can see that the bands are spin splitting at the $K (K' \prime)$ valley when both $\lambda_E$ and $\lambda_{AFM}$ are finite. More importantly, the exchange splitting can be switched by changing the direction of electric field $E_z$ through dual-gate tunability. We note that the theoretical results in figures 1(b) and (d) are obtained by using the properties $E_{\varepsilon,s}(\lambda_E) = E_{\varepsilon,s}(-\lambda_E)$ in equation (3) directly. Such electric field effects on magnetism in BLG are, in general, attributed to spin splitting that supports spin polarization of flowing electrons along certain directions.

As commented in the introduction, we propose and study a vdW-based MR device. The geometry is shown in figure 2(a). The device consists of two semi-infinite electrodes connected by a BLG as the channel material of the central scattering region. The left electrode and right electrode are vdW spin valves. The proximity effect-induced magnetism in the right vdW spin valve is electrically tunable by the dual-gated structure. The device is analyzed by $P (E_z \downarrow)$ and $AP (E_z \uparrow)$ magnetization alignments of the left and right electrode.

We describe each part of the device with the 2D low-energy Hamiltonian for the two-component wave function $\Psi = (\psi_{A1}, \psi_{B2})^T$ (subscript $A, B$ refers to sublattices in each layer of graphene, while subscript 1, 2 represents the two layers of graphene), which gives an excellent description of the electronic properties of...
each region of the device near the momentum valleys, is expressed as [45, 56, 64, 65]

\[ H_{\sigma i} = H_0 + \lambda_E (x) \sigma_z - s \lambda_{AFM} (x) \sigma_0 + s \lambda_{AFM} (x) \sigma_z - \mu (x) \sigma_0, \]

with

\[ \lambda_E (x) = 2F_{E_{F}} \Theta (x - L), \quad \lambda_{AFM} (x) = h \Theta (-x), \]

\[ \lambda_{AFM} (x) = \Delta \Theta (x - L), \quad \mu (x) = U \Theta (x) \Theta (L - x), \]

Here, \( \lambda_{AFM} (x) \) with strength \( h \) is the spin splitting in the left vdW spin valve with the P magnetizations of the two IFMs. \( \mu (x) \) is a strength \( U \) is an electrostatic potential due to the doping [66–69]. Note that \( \sigma_0 \) is a \( 2 \times 2 \) identity matrix, and \( \Theta (x) \) is the Heaviside step function to characterize the local existence of the spin splitting, the band gap, and the electric field \( E_z \).

According to the equation (4), the band structures around the momentum valleys in the three different regions are

\[ E_{s,i,j} = \pm \sqrt{\left( \frac{\hbar^2 k^2}{2m} \right)^2 + (\lambda_E + s\Delta)^2 \Theta (x - L) - s\mu \Theta (-x),} \]

+ \( U \Theta (x) \Theta (L - x). \]

According to equation (6), the corresponding band alignments in different spatial regions are plotted in figures 2(b) and (c) for potential profiles of barrier \( (n - p - n) \) and well \( (n - n' - n) \), respectively. Obviously, \( \lambda_{AFM} \) at \( x < 0 \) splits the spectrum into two branches with breaking spin degeneracy, while imposing the pure perpendicular electric-field modulation \( E_z \) at \( x > L \), the combination of \( \lambda_E \) and \( \lambda_{AFM} \) modulate the band structure and explicitly break the spin degeneracy. In addition, the low-energy spectrum of BLG for \( 0 < x < L \) is gapless and spin-degenerate.

We now turn our attention to the corresponding eigenfunctions in each region, which can be obtained from the equation \( H_{s,i} \psi_{s,i} = E \psi_{s,i} \) with Hamiltonian \( H_{s,i} \) given in equation (4), have the forms,

\[ \psi_{R,s,i} (x) = \frac{1}{N_{R,s}} \left( \begin{array}{c} \delta_{+s,i} \\ \delta_{-s,i} \end{array} \right) \exp \left( ik_{s,i} x + ik_{y,i} y \right), \]

\[ \psi_{L,s,i} (x) = \frac{1}{N_{L,s}} \left( \begin{array}{c} \delta_{+s,i} \\ \delta_{-s,i} \end{array} \right) \exp \left( -ik_{s,i} x + ik_{y,i} y \right), \]

and

\[ \phi_{R,s,i} (x) = \frac{1}{\gamma_{s,i}} \left( \begin{array}{c} \frac{\delta_{+s,i}}{N_{R,s}} \\ \frac{\delta_{-s,i}}{N_{R,s}} \end{array} \right) \exp \left( -\kappa_{s,i} x + ik_{y,i} y \right), \]

\[ \phi_{L,s,i} (x) = \frac{1}{\gamma_{s,i}} \left( \begin{array}{c} \frac{\delta_{+s,i}}{N_{L,s}} \\ \frac{\delta_{-s,i}}{N_{L,s}} \end{array} \right) \exp \left( \kappa_{s,i} x + ik_{y,i} y \right), \]

where \( \psi_{R(s),s,i} (x) \) is the propagating eigenmode along the positive (negative) \( x \)-axis with wavevector \( k_{s,i} (x) = \sqrt{2T_{s,i} (x) - k_{y,i}^2}. \) Here, \( T_{s,i} (x) = \sqrt{(E + s\lambda_{AFM} (x) + \mu (x))^2 - (s\lambda_{AFM} (x) + \lambda_E (x))^2}. \) \( N_{R(s),s,i} (x) \) describes not the propagation but the exponential attenuation of the amplitude with increasing \( x \), which contains a damping parameter \( \kappa_{s,i} (x) = \sqrt{2T_{s,i} (x) - k_{y,i}^2}. \) \( N_{R(s),s,i} = \sqrt{(k_{s,i}^2 + k_{y,i}^2)^2 + \delta_{2,i}^2} \) is the normalized constant with \( \delta_{2,i} = 2 [E + \mu (x)] \pm 2 [s\lambda_{AFM} (x) + \lambda_E (x)] \pm 2s\lambda_{AFM} (x) \). Also note that \( \pi_{s,i} = i\hbar k_{s,i} \pm k_{y,i} \) and \( \gamma_{s,i} = k_{s,i} - \eta n_{s,i} \) where \( k_{s,i} \) is the wave vector parallel to \( y \)-axis.

In the absence of valley-symmetry-breaking disorders (leading to valley mixing), the valley degeneracy and valley symmetry are preserved, and hence we below consider incident electrons in a single valley with the incident energy being \( E_F \), spin \( s (s = \pm 1) \) is the eigenvalue of the operator \( s_z \), momentum \( (k_{s,i}, k_{y,i}) \) and incident angle \( \alpha \) with respect to the \( k_x \) axis, the wave functions \( \Psi_L, \Psi_C, \Psi_R, \) regarding \( L \) and \( R \) as semi-infinite leads and modeling \( C \) as the scattering region, are given by

\[ \Psi_L (x) = \psi_{R,s,i} (x < 0) + r_{s,i} \psi_{L,s,i} (x < 0), \]

+ \( \cos \alpha \psi_{L,s,i} (x < 0) \),

\[ \Psi_C (x) = a_1 \psi_{R,s,i} (0 < x < L) + a_2 \psi_{L,s,i} (0 < x < L), \]

+ \( a_3 \phi_{L,s,i} (0 < x < L) + a_4 \phi_{L,s,i} (0 < x < L), \)

\[ \Psi_R (x) = \psi_{R,s,i} (x > L) + d_{s,i} \phi_{L,s,i} (x > L), \]

We can uniquely determine the coefficients \( r_{s,i}, \cos \alpha, a_1, a_2, a_3, a_4 (m = 1–4) \) in equation (9) by imposing the continuity of the scattering state and the continuity of the first derivative at \( x = 0 \) and \( x = L \) [70], which can be written as
respectively. The eigenvectors in figure 3, seeing the caption for details on the parameters. First, as depicted in the figures 3(a)–(d), we following [20],

where the velocity operator \( \tilde{v}_x \) in the \( x \)-direction therefore reads:

\[
\tilde{v}_x = \frac{\partial H}{\hbar \partial k_x} = \begin{pmatrix} 0 & i \frac{\partial}{\partial x} \\ -i \frac{\partial}{\partial x} & 0 \end{pmatrix} + \begin{pmatrix} 0 & \imath \eta k_y \\ -\imath \eta k_y & 0 \end{pmatrix}
\]

Note that \( \tilde{v}_x \) is a \( \eta \)-dependent velocity operator, while the group velocity \( v_{RL}(x) = \langle \psi_{R}(x)|i\partial \psi_{L}(x)/\partial t \rangle \) is independent of the valley index \( \eta \). This fixes the spin and valley-resolved transmission probabilities \( T_{Ez}^{Ez}(E_{z}) = \frac{\langle \tilde{\psi}_{E_{z}^{\dagger}}(E_{z}) | \psi_{E_{z}}(E_{z}) \rangle^{2}}{\langle \tilde{\psi}_{E_{z}^{\dagger}}(E_{z}) | \tilde{\psi}_{E_{z}^{\dagger}}(E_{z}) \rangle} \) in all the outgoing propagating modes for the negative \( (E_{z}, \downarrow) \) and positive \( (E_{z}, \uparrow) \) electric-field configurations. We then obtain the 2D two-terminal zero temperature conductance in the device by using

\[
G_{E_{z}^{\uparrow}} = G_{0} \sum_{s} \int_{-\frac{T}{2}}^{\frac{T}{2}} P_{s} T_{E_{z}^{\uparrow}}^{E_{z}^{\dagger}}(E_{z}) \cos \alpha \ d\alpha
\]

with

\[
T_{E_{z}^{\uparrow}}^{E_{z}^{\dagger}}(E_{z}) = T_{E_{z}^{\uparrow}}^{E_{z}^{\dagger}}(E_{z}) + T_{E_{z}^{\uparrow}}^{E_{z}^{\dagger}}(E_{z})
\]

where \( G_{0} = \frac{e^{2}}{h} \) and \( P_{s} = \frac{1}{2} \left( 1 + s \frac{\hbar}{\omega_{q}} \right) \). With these definitions, we can introduce the corresponding MR as following [20],

\[
MR = \frac{G_{E_{z}^{\uparrow}} - G_{E_{z}^{\dagger}}}{G_{E_{z}^{\dagger}}}
\]

3. Numerical results and discussions

We shall see that the main ingredients to compute the MR effect are the spin-resolved transmissions. Thus, the question we must address now is the determination of these values for the structure studied in this work. Before showing the numerical results, it is instructive to give the following comment. Without losing generality, throughout this work we will concentrate only on the \( n - p - n \) junction and the \( n - n'-n \) junction, as the results of the \( n - p - n \) junction and the \( p - p' - p \) junction can be obtained by the symmetry of the effective Hamiltonian [equation (4)], which is given by

\[
C_{x} C_{y} H_{xy} (k_{x}, k_{y}) C_{y} C_{x} = H_{-xy} (k_{x}, k_{y}) | \mu(x) \rightarrow -\mu(x), \lambda_{x}(x) \rightarrow -\lambda_{x}(x), \rangle
\]

where \( C_{x(y)} \) and ‘*’ are the inversion of the spin(valley) state \( s \rightarrow -s \) (\( \eta \rightarrow -\eta \)) and the complex conjugate, respectively. The eigenvectors \( \tilde{\psi}_{E_{z}^{\dagger}} \) and \( \phi_{E_{z}^{\dagger}} \) of equation (4) are transformed as

\[
\begin{align*}
C_{x} C_{y} \tilde{\psi}_{E_{z}^{\dagger}}' (k_{x}, k_{y}) & = \tilde{\psi}_{-E_{z}^{\dagger}} (k_{x}, k_{y}) | \mu(x) \rightarrow -\mu(x), \lambda_{x}(x) \rightarrow -\lambda_{x}(x), \rangle \\
C_{x} C_{y} \phi_{E_{z}^{\dagger}}' (\kappa_{x}, k_{y}) & = \phi_{-E_{z}^{\dagger}} (\kappa_{x}, k_{y}) | \mu(x) \rightarrow -\mu(x), \lambda_{x}(x) \rightarrow -\lambda_{x}(x), \rangle
\end{align*}
\]

Applying equations (9)–(11), we find that the transmission probability for the \( n - p - n \) \( (n - n' - n) \) junction and the \( p - n - p \) \( (p - p' - p) \) junction are identical, i.e.

\[
T_{E_{z}^{\uparrow}}^{E_{z}^{\dagger}}(E_{z}) = T_{E_{z}^{\dagger}}^{E_{z}^{\uparrow}}(E_{z}) = T_{E_{z}^{\dagger}}^{E_{z}^{\uparrow}}(E_{z})
\]

It follows that the charge conductance, which is obtained by using equation (12), for the \( n - p - n \) \( (n - n' - n) \) junction in the negative \( (E_{z}, \downarrow) \) electric-field configurations equals that for the \( p - n - p \) \( (p - p' - p) \) junction in the positive \( (E_{z}, \uparrow) \) electric-field configurations, \( G_{E_{z}^{\uparrow}}^{p-p-n} = G_{E_{z}^{\dagger}}^{p-p'-p} \). Overall, based on equation (14), we find the following relation is satisfied:

\[
MR_{p-p-n(a-b)} = \frac{MR_{p-p'-p(a-b)}}{2}
\]

We now focus on the contour plots of spin-resolved transmissions \( T_{E_{z}^{\dagger}}^{E_{z}^{\uparrow}}(E_{z}) \) as a function of incident angle \( \alpha \) and electric potential \( U \), which are calculated numerically by means of equation (13) and presented in figure 3, seeing the caption for details on the parameters. First, as depicted in the figures 3(a)–(d), we show \( T_{E_{z}^{\dagger}}^{E_{z}^{\uparrow}}(E_{z}) \) for the \( n - p - n \) junction [see figure 2(b)]. At certain non-normal incidence, one can clearly
see that increasing the height of the potential barrier leads to FP resonances. The shape of the narrow resonances in the \((\alpha, U)\) plane can be understood by applying the condition \(k_{\alpha}L = n\pi\), which is due to an interference of Dirac electron in the scattering region. For angles close to \(\alpha = 0\), however, the transmissions for normal incidence is practically zero regardless of the barrier height; i.e. the Klein tunneling with unit probability occurring in monolayer graphene is replaced by the suppression of normal transmissions, which is called anti-Klein tunneling [35, 37, 71]. This perfect reflection for electrons traversing the barrier with normal incidence is the consequence of the pseudospin polarization vectors of the massive Dirac electron in \(n\) and \(p\) regions are orthogonal to each other [42, 71]. Incidentally, we also observe that when moving away from normal incidence, the transmissions in spin-up and spin-down channels increase from zero and, after the first critical angle the transmissions are almost fully suppressed. Compare figures 3(a) and (c) with figures 3(b) and (d), we also notice that the spin-resolved transmissions in the \(E_z \downarrow\) and \(E_z \uparrow\) configuration are different. At a given height of the potential barrier, such as \(\frac{U}{\Delta} = 9, 14\), or \(-17\), we find that: (1) the transmission in spin up channel for \(E_z \downarrow\) configuration is almost fully suppressed while that for the \(E_z \uparrow\) configuration is large, i.e. \(T^{E_z\downarrow}_{\alpha}(\alpha) \gg T^{E_z\uparrow}_{\alpha}(\alpha) \approx 0\), as shown in figures 3(a) and (b); (2) the results in figures 3(c) and (d), however, demonstrate that the spin-down transmission shows almost opposite behavior, \(T^{E_z\uparrow}_{\alpha}(\alpha) \gg T^{E_z\downarrow}_{\alpha}(\alpha) \approx 0\).

Next, we consider the \(n - n' - n\) case [see figure 2(c)], as shown in figures 3(e)–(h). In this case, the oscillatory behaviors of the transmissions for both negative \(E_z \downarrow\) and positive \(E_z \uparrow\) configurations become somewhat more pronounced but remain quite similar qualitative features to potential barrier. As seen, similar to the case of potential barrier, for certain angles, the transmissions are also periodic as the depth of potential well increases. However, instead of the perfectly reflected for angles close to \(\alpha = 0\) [see figures 3(a)–(d)], massive Dirac electron injecting normal to the junction is not perfectly reflected, i.e. high transmissions are found at normal incidence, which is apparent in figures 3(e)–(h). This finite transmissions occur because the pseudospins in the monopolar \(n - n' - n\) junction are parallel [42, 63].

We further analyze the impact of the \(E_z \downarrow\) and \(E_z \uparrow\) configurations on the spin-resolved transmissions. As shown in figures 3(e) and (f), for \(\frac{U}{\Delta} = 7.9\), or 15 with \(\alpha\) in the transmission window, \(T^{E_z\downarrow}_{\alpha}(\alpha) \gg T^{E_z\uparrow}_{\alpha}(\alpha) \gg 1.0\) while in figures 3(g) and (h), \(T^{E_z\downarrow}_{\alpha}(\alpha) \gg T^{E_z\uparrow}_{\alpha}(\alpha) \gg 0.5\). Additionally, we obtain that the carrier transmissions increase significantly for the incident angle within the allowed transmission window. Especially, the broadening of the resonance peaks enhance transmissions for angles in the vicinity of the resonance condition. The distinct discrepancy of the spin-resolved transmission between the barrier \((n - p - n\) junction) and the well \((n - n' - n\) junction) manifest themself by the presence of many more and well-defined resonances and a substantially higher transmissions for \(n - n' - n\) junction. We will see
that the resulting conductance of the \( n - p - n \) junction incorporating the anti-Klein tunneling differs significantly from the conductance calculated within the \( n - n' - n \) junction.

Let us now move on to look at the impact of the height (depth) of potential barrier (well) \( U \) on the performance of MR and configuration-resolved conductance \( G_{E_z \downarrow} / G_0 \) and \( G_{E_z \uparrow} / G_0 \), as shown in figure 4. One can see from figure 4(a) that the MR show nice and regular damped oscillations with the height of potential barrier \( U \). Furthermore, increasing the strength of \( \lambda_E \) can even produce a giant MR effect with magnitude up to 1200\%, which is almost seven times larger than that in the graphene-based FM double junctions [72].

Interestingly, changing the height of potential barrier \( U \) value leads to transitions from positive MR to small negative MR, which means that the spin filtering property can easily be switched by tuning \( U \). Those can be understood from the configuration dependence of conductance in figures 4(b) and (c). It is clear to see that the oscillatory behavior of the MR with \( U \) originate from \( G_{E_z \downarrow} \) and \( G_{E_z \uparrow} \), which all exhibit oscillate behaviors. This can also be explained in the context of the FP oscillations of transmissions [see figures 3(a)–(d)]. Also, we can see that the opposite shifts of the conductance peaks for \( E_z \downarrow \) with respect to those for \( E_z \uparrow \) can provide a negatively stable MR, which mainly comes from the different transport properties of up (down) spin-Dirac electron in the positive and negative values of \( E_z \).

In figure 5, we present the dependence of MR (panels (a) and (d)), \( G_{E_z \downarrow} \) (panels (b) and (e)) and \( G_{E_z \uparrow} \) (panels (c) and (f)) on the width \( L \) of the scattering region for \( \lambda_{E_z \downarrow} = -\lambda_{E_z \uparrow} = -3\Delta \), \( \lambda_{E_z \downarrow} = -\lambda_{E_z \uparrow} = -2\Delta \) and \( \lambda_{E_z \downarrow} = -\lambda_{E_z \uparrow} = -\Delta \), respectively. Also two values of the electrostatic potential are considered: \( U = -4\Delta \) (left side panels) and \( U = 4\Delta \) (right side panels). In the case of the potential barrier \( U = -4\Delta \), MR exhibits a nonmonotonic behavior with a maximum at around \( k_FL = 6 \). One can note that when \( \lambda_{E_z \downarrow} = -\lambda_{E_z \uparrow} = -2\Delta \) is present, the value of MR becomes larger and even reaches 1000\%. We emphasize that a finite potential barrier \( U = -4\Delta \) makes negative MR effect and can even result in a giant MR effect with a sign oscillating with the barrier width \( L \), which also produces an opposite shift of the conductance peaks from the negative \( (E_z \downarrow) \) to the positive \( (E_z \uparrow) \) electric-field configurations, as illustrated.
in figures 5(b) and (c). Contrary to the case presented in figure 5(a), for the potential well \((U = 4\Delta)\), MR is positive and could not rise above 200\% [see figure 5(d)]. When comparing figures 5(e) and (f), one can note that the conductance is larger in the negative \((E_z \downarrow)\) configuration than that in the positive \((E_z \uparrow)\) one. Moreover, for both types of junctions, the oscillations depend on the quantum interference behavior of different propagating wave vectors for two different spin indices in the scattering region.

Finally, a comment is in order on the feasibility of the structure we considered. An essential ingredient for our structure is vdW spin valves based on BLG, in which a perpendicular electric field can be used to control the spin splitting. According to DFT calculations [45, 73, 74], the magnetic proximity-induced spin orbit interactions (SOI) in the vdW spin valve structure is expected to be weaker than the spin proximity effect discussed here, at this stage, the spin flip caused by the SOI is not taken into account in our calculation. Besides, we can expect the origin of the theoretically calculated results to be mainly caused by the simple quadratic relation due to the absence of trigonal warping effect by assuming \(|E - U| > 0.005\gamma_1\) [42, 63]. Furthermore, the electrostatic potential \(U\) in scattering regions can be adjusted not only by doping mentioned in this paper, but also by gate voltages [54-57]. From similar studies, we think that our junctions can be realized experimentally. Thus, our results suggest an alternative way to create and electrically tunable MR, which can be used as lower power magneto-electric devices.

4. Conclusion

In summary, we propose a design of a MR device based on 2D BLG, in which two vdW spin valves are used as the FM leads. The vdW spin valve enables modulation of the direction and magnitude of the magnetization by out-of-plane electrical field. We show that MR and conductance all exhibit oscillate behaviors due to the quantum interference of propagating waves for two different spin indices in the scattering region. In the case of \(n - p - n\) junction, MR becomes negative for certain values of the height \(U\) of the barrier or the width \(L\) of the scattering region. This is because the opposite shifts of the conductance peaks for \(E_z \downarrow\) with respect to those for \(E_z \uparrow\). Moreover, a giant MR effect can be created with increasing
λ_Ez. On the other hand, for the n′ − n junction, the MR oscillations are also regular but the amplitude of the MR signal is positive and much smaller than that in the previous case. This is due to a different role of GEz and GEz0, which have the same oscillation periods but have the different magnitude. Such an electrically tunable MR device can open up new possibilities for lower power magneto-electric devices since no current-induced magnetization switching.

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Data availability statement

All data that support the findings of this study are included within the article (and any supplementary files).

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