Electric-field-induced extremely large change in resistance in graphene ferromagnets

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Abstract
A colossal magnetoresistance (∼100 × 10^3%) and an extremely large magnetoresistance (∼1 × 10^6%) have been previously explored in manganite perovskites and Dirac materials, respectively. However, the requirement of an extremely strong magnetic field (and an extremely low temperature) makes them not applicable for realistic devices. In this work, we propose a device that can generate even larger changes in resistance in a zero-magnetic field and at a high temperature. The device is composed of graphene under two strips of yttrium iron garnet (YIG), where two gate voltages are applied to cancel the heavy charge doping in the YIG-induced half-metallic ferromagnets. By calculations using the Landauer–Büttiker formalism, we demonstrate that, when a proper gate voltage is applied on the free ferromagnet, changes in resistance up to 305 × 10^6% (16 × 10^3%) can be achieved at the liquid helium (nitrogen) temperature and in a zero magnetic field. We attribute such a remarkable effect to a gate-induced full-polarization reversal in the free ferromagnet, which results in a metal-state to insulator-state transition in the device. We also find that the proposed effect can be realized in devices using other magnetic insulators, such as EuO and EuS. Our work should be helpful for developing a realistic switching device that is energy saving and CMOS-technology compatible.

Keywords: graphene ferromagnets, extremely large change in resistance, electric-field-induced full-polarization reversal, magnetic proximity effect, magnetic insulators

(Some figures may appear in colour only in the online journal)
on-off ratios. A strong magnetic field is required for a relatively small low-resistance.

Several years ago, a colossal negative magnetoresistance was explored in functionalized graphene such as doped fluorinated graphene [8, 9]. Very recently, an extremely large magnetoresistance (XMR) was explored in other Dirac materials. For example, unsaturated XMR up to $0.45 \times 10^6$ at 4.5 K in a magnetic field of 14.7 T and XMR up to $13 \times 10^6$ at 0.53 K in a magnetic field of 60 T were observed in WTe$_2$ [10]. XMR of about one million percent at 4.2 K and $9 T$ was observed in LaSb [11], XMR of $0.1 \times 10^6$ and $0.73 \times 10^6$ at 2.5 K and 14 T were obtained in NbAs$_2$ and TaAs$_2$, respectively [12], and unsaturated XMR up to $11.2 \times 10^6$ at 1.8 K in a magnetic field of 33 T was obtained in PtBi$_2$ [13]. However, to achieve these CMR or XMRs, extreme high-magnetic fields and low-temperatures are required. In these cases, it is not applicable for any realistic devices yet.

In this work, we propose a device that can generate even higher on-off ratios at a high temperature and in a zero magnetic field. The device is composed of graphene under two strips of a magnetic insulator yttrium iron garnet (YIG), where two gate voltages are applied to cancel the heavy charge dopings in the YIG-induced half-metallic ferromagnets (see figure 1(a)).

Based on calculations using the Landauer–Büttiker formalism, we demonstrate that, by applying a proper voltage on one of the ferromagnets (the free one), an extremely large change in resistance up to $305 \times 10^6$ can be achieved at the liquid helium temperature (4.2 K) and in a magnetic filed of 0 T. The value maintains as $16 \times 10^6$ at the liquid nitrogen temperature (77 K). We indicate that such huge on-off ratios stem from an electric-field-induced reversal of the full polarization in the free ferromagnet, which results in a transition between a metal and an insulator state in the device. Moreover, we find that the proposed effect can be realized in graphene under two strips of other magnetic insulators, such as EuO and EuS.

2. Device setup

The designed device is shown in figure 1(a). A sufficiently wide (in the y-direction) and short (in the x-direction) graphene strip of $L \times W$ is grown on a substrate. Here, $W$ is several times that of $L$ to ensure that the transport is dominated by bulk states and does not depend on the edge types [14]. On top of the graphene, two YIG strips of lengths $l_p$ and distance $d$ are transferred [15]. As demonstrated by experiments [15–18] and first-principle calculations [19], the YIG strips induce ferromagnetism in the underlying graphene through an overlap of the spin-polarized Fe-3d states and the C $p_z$ states, which is called a magnetic proximity effect. The electronic structure of the graphene ferromagnet can be described by an exchange splitting for electrons or holes ($\delta_e$ or $\delta_h$) and a band gap opening at the Dirac point ($E_D$), see figure 2(b). Unfortunately, a heavy electron doping ($E_D = -0.78$ eV) is also induced in the graphene ferromagnet [15, 19], which limits its spintronic application by a low polarization.

On the YIG strips, the two top gates ($V_{p,f}$) are placed [15]. They are used to cancel the heavy dopings through a strong electric-field effect [15] similar to that in the pristine graphene [20]. The substrate under graphene contacts a back gate ($V_b$), which is used to tune the Fermi energy ($E_F$)
through the whole device. When the Fermi energy is set as $E_F \in (E_G/2, E_G/2 + \delta_e)$ or $E_F > E_G/2 + \delta_e$, a half-metallic or normal ferromagnet of electron polarity is made use of. A small voltage ($\Delta V_J$) is further applied on the free ferromagnet to tune its Dirac point ($E_{DJ}$). As we will see below, the Fermi energy and the Dirac point of the free ferromagnet are important to the states of the device. On the two sides of the device, the graphene film contacts with a source and drain electrode, which may induce a contact doping ($U$) and a contact resistance [21–23]. Instead of YIG, EuO and EuS can be deposited on the graphene film. The ferromagnetism is induced by an overlap of the Eu – 4f states and the C – p$_c$ states. All of the parameters ($\delta_{eh}^{(i)}$, $E_{b}^{(i)}$, $E_{G}^{(i)}$) [19] as well as the Curie temperatures ($T_c^{(i)}$) [15, 24, 25] become different, see figures 2 (a) and (b). Here, $i = 1,2,3$ for YIG, EuO and EuS, respectively.

### 3. Calculation formula

Below, we will derive the formula to calculate the device conductance at different $E_F$ and $E_{DJ}$. Devices using YIG, EuO and EuS will be handled uniformly.

The low-energy dispersions around the Dirac points, which are predicted by first-principle calculations [19, 26], are recalcualted and plotted in figure 2. The dispersions read

$$E_i^{(i)}(q) - E_{D}^{(i)} = D_i^{(i)} \pm \sqrt{(h\nu_i q)^2 + (\Delta_i^{(i)})^2},$$

where $i = 1,2,3$ are the material- and spin-dependent Dirac cone dopings, Fermi velocities and Dirac gaps, respectively. From figure 2, one can see clear spin-resolved parabolic dispersions, which accompany the band gap opening at the Dirac point and exchange splittings for the electrons and holes. The parameters for graphene on six trilayer YIG (six bilayer EuO and EuS) are given in [19]. The parameters in equation (1) relate to them by $D_i^{(i)} = \pm \delta_{eh}^{(i)}/2$ and $\Delta_i^{(i)} = |\delta_{eh}^{(i)}| + E_{G}^{(i)}$, see figure 2(b); $\nu_1^{(i)}$ are fitted from the original dispersions [19]. The values are listed in table 1.

Regarding equation (1) as a combination of dispersions of two gapped graphene, the dispersions of the ferromagnets can be described by an uniform effective Hamiltonian in a sublattice space [27]

$$H_{k,k'}^{(i)} = \mathcal{I}(D_i^{(i)} + E_{D}^{(i)}) + \xi \sigma_i \Delta_i^{(i)} + \sigma_i h\nu_i^{(i)} k,$$

where $\mathcal{I}$ is the identity matrix, $\xi = \pm 1$ is the index for valley $K$ and $K'$, $\sigma = (\sigma_x, \sigma_y)$ is the pseudospin Pauli matrices and $k = (k, q)$ is the momentum operator. Such an effective Hamiltonian is different from those described in a sublattice-spin direct produced space [19, 28, 29], where the ferromagnets are viewed as a Dirac gap and two exchange splittings. Note, in this space, a four-component wave function should be solved [30]. Different from them, in equation (2), we also consider a spin-resolved Fermi velocity $v_i$. As we will show below, this parameter determines an effective spin Dirac gap ($\Delta_i^{(i)}$), hence playing an important role in the insulator state. For the graphene under the electrode metals, $\mathcal{H}_{k,k'} = \mathcal{I}U + \sigma \cdot h\nu_{k,k'} k$, and for the pristine graphene between the pinned and free ferromagnets, $\mathcal{H}_{k,k'} = \mathcal{I}(\phi(x) + \sigma \cdot h\nu_{k})$, where $\phi(x) = (x/d)E_{D}^{(i)}$ is the potential shift.

When $E_{D}^{(i)} = 0$ (as is the case in figures 1(b1) and (c1)), the right- and left-going envelope functions in the contacted, ferromagnetic and pristine graphene ($i = c,m,p$) can be obtained by exactly resolved $H_{i}\mathcal{F}_{i}^{(i)} = E_{i}\mathcal{F}_{i}^{(i)}$. Using a characteristic energy (length) $E_0 = 10$ meV ($l_0 = h\nu_F/E_0 = 56.55$ nm), the dimensionless result reads

$$\Phi_{i}^{(i)} = \frac{1}{\sqrt{2E_0}} \left( \frac{E_j}{(\pm k_j + iq_j)} \right) e^{\pm ik_x^{(i)}+iq_y^{(i)}},$$

where $E_{p(c)} = E(-U)$, $E_m = (E_s + \Delta_s)/v_s$ with $E_s = E - D_s(-E_{D}^{(i)})$ for the pinned (free) ferromagnet, $q_{p,c,m} = E_s \pm \alpha$, $\alpha$ is the conserved transverse wave vector, $k_{p,c} = \text{sign}(E_{p,c}) (E_{p,c} - q_{p,c}^2)^{1/2}$ and $k_{m} = \text{sign}(E_m)(E_m - q_m^2)^{1/2}$ with $E_m = (E_s - \Delta_s)/v_s$. It is seen that $k_{s}^{2} + q_{s}^{2} = (E_s^{2} - \Delta_s^{2})/v_s^{2}$. This means that an effective spin Dirac gap, $\Delta_s/v_s$, is determined by not only the spin-dependent Dirac gap but also the spin-dependent Fermi velocity.

When $E_{D}^{(i)} > 0$ (the cases in figures 1(b2) and (c2)), the graphene between the ferromagnets becomes an n-p junction and the envelope function cannot be straightforwardly solved [31, 32]. Instead, the function can be solved in a pseudospin space, which is rotated by $\pi/2$ around the y-axis (see figure 1(a)) [31, 32]. The result reads

$$\Phi_{p} = e^{\mp \left( \frac{F(x)}{G^{*}(x)} \right) e^{i\theta_{p}} + \frac{e}{c^{+}} \left( \frac{G(x)}{F^{*}(x)} \right) e^{i\theta_{p}}},$$

where $F(x) = D[-1 + iq^2/2f, (1 + i)(E + fx)/\sqrt{f}]$ and $G(x) = (1 + i)\sqrt{f}q^{-1}D[|q^2|/2f, (1 + i)(E + fx)/\sqrt{f}]$ with $D[.]$ being the Weber parabolic cylinder function and $f = E_{D}/d$. Note that $F(x)$ and $G(x)$ ($F^{*}(x)$ and $G^{*}(x)$) have the properties of a right (left)-going evanescent wave function [31]. In the rotated pseudospin space, the envelope functions for the contacted and ferromagnetic graphene become $\Phi_{i}^{(i)} = [E_j \pm k_j + iq_j, -E_j \pm k_j + iq_j]^{T}(2E_{i}^{(i)})^{-1}e^{\pm ik_x^{(i)}+iq_y^{(i)}}$.

For $T < \min(100$ K, $\frac{1}{2}$), the ferromagnetics hold and the inelastic scatterings (e-e and e-ph) can be ignored [33, 34]. The spin-resolved conductance through the device can be given by the Landauer–Büttiker formula [35]

$$G_{i}^{(i)}(E_{D}, E_F, T) = G_0 \int dE \frac{-d\Gamma(E, T)}{dE} \int_{-\pi/2}^{\pi/2} \left| \mathcal{F}_{i}^{(i)}(E, T) \right|^2 \cos \alpha_d \alpha_d,$$

### Table 1. The Curie temperatures, Dirac cone dopings (in the unit of eV), spin Dirac dopings (in the unit of $E_0$), spin Dirac gaps (in the unit of $E_0$) and spin Fermi velocities (in unit of $v_F$) for graphene under YIG, EuO and EuS.

| $T_{c}$ (K) | $E_{D}$ | $D_{i}$ | $\Delta_{i}$ | $\nu_{i}$ | $\Delta_{i}$ | $\nu_{i}$ |
|------------|---------|--------|-------------|---------|-------------|---------|
| YIG/Gr     | 550     | 0.78   | 5.8         | 11.6    | 0.63        | 2.6     | 5.3         | 0.70    |
| EuO/Gr     | 69      | 1.37   | 4.2         | 13.4    | 1.34        | -2.4    | 9.8         | 1.63    |
| EuS/Gr     | 16.5    | 1.5    | 1.6         | 18.3    | 1.40        | 0       | 15          | 1.70    |
and with a standard transfer matrix method [36]. The enters into the orange window. This leads to a huge on-off ratio of 250002. (a), (c), (e) The cases for devices with half-metallic ferromagnets using YIG, EuO, and EuS (EF = 2.5, 7, 9.5), respectively. The smaller (larger) spin component is reduced (enlarged) by 2 times for clearness. (b), (d), (f) The cases for devices with normal ferromagnets using YIG, EuO, and EuS (EF = 15), respectively.

where \( G_0 = 2e^2/h \times (|E_F|W/h\nu_F) \) is the unit conductance, \( f(E, T) = [1 + e^{(E - E_F)/kT}]^{-1} \) is the Fermi–Dirac distribution function and \( t_s(E_{DF}, E, \alpha) \) is the spin-resolved transmission coefficient at a Dirac point of \( E_{DF} \), an energy of \( E \) and an incident angle of \( \alpha \). \( t_s^{(i)}(E_{DF}, E, \alpha) \) can be solved by \( \Phi_j \) with a standard transfer matrix method [36]. The \( E_{DF} \)-dependent resistance of the device can be given by 
\[
R(E_{DF}, E_F, T) = [G^{(i)}(E_{DF}, E_F, T) + G^{(f)}(E_{DF}, E_F, T)]^{-1} \times \text{unit conductance},
\]
where \( R_0 = G_0^{-1} \). The changes in resistance (RC) at different \( E_{DF} \) can be defined as
\[
RC^{(i)}(E_{DF}, E_F, T) = \frac{R_0(E_{DF}, E_F, T) - R_0(M, 0, E_F, T)}{R_0(M, 0, E_F, T)} \times 100\%.
\]

To calculate a change in resistance, four conductance \( G^{(i)}(0, E_F, T) \) and \( G^{(f)}(E_{DF}, E_F, T) \), should be calculated. Each of them depends strongly on the magnetic insulator (\( l_p \)), ferromagnet length (\( l_p \)), electrode contact (\( U \)), Fermi energy (\( E_F \)) and temperature (\( T \)).

4. Results and discussion

4.1. Gate-induced huge changes in resistance based on metal-insulator states

We first consider a YIG-based device with two half-metallic ferromagnets (\( E_F = 2.5 \)) of \( l = 1 \). Figure 3(a) shows the zero-temperature device resistance as a function of \( E_{DF} \). As can be seen, the device resistance is rather low (\( 10^4 \) \( R_0 \)) at \( E_{DF} = 0 \); it becomes rather high (\( 10^8 \) \( R_0 \)) as \( E_{DF} \) enters into the orange window. This leads to a huge on-off ratio of 3.92 \( \times 10^8 \%) at \( E_{DF} = 3.85 \). At \( E_{DF} = 0 \), both ferromagnets are half-metallic (full polarizations of spin-up). Accordingly, electrons of spin up from the device source are transparent, while electrons of spin down are blocked by two potential barriers (see figure 1(b)). These behaviors result in a near-unit transmission for spin up and a rather small transmission (\( 10^{-6} \)) for spin down, respectively, see the dashed curves in figure 4(a). The transmissions respectively contribute a rather small (0.62 \( R_0 \)) and large (\( 1.5 \times 10^7 \) \( R_0 \)) resistance, see the dashed curves in figure 3(a). The total resistance, as a parallel connection of them, is rather small (0.62 \( R_0 \) almost equal to the smaller one). On the other hand, as can be seen in figure 5(a), the low resistance increases with an increasing temperature, which implies that the device supports a metal state at \( E_{DF} = 0 \).

In the orange window, \( E_{DF} \in (3.0, 13.8) \) and \( E_{DF} - E_{DF} \in (11.3, -0.5) \in (-|\delta_n| - E_G/2, -E_G/2) \), see figure 2. The latter means that the Fermi energy lies in the hole exchange splitting window of the free ferromagnet, see figure 1(b). Such a reversal of the charge polarity leads to an reversal of the spin polarization. In other words, by a proper \( E_{DF} \), the spin-up full polarization in the free ferromagnet becomes a spin-down full polarization. As a result, electrons of spin up from the device source, which are transparent at \( E_{DF} = 0 \), now encounter a barrier in the free ferromagnet. On the other hand, electrons of spin down from the device source, which were blocked by two barriers \( E_{DF} = 0 \), are now still blocked by a barrier in the pinned ferromagnet. In the orange window, where \( E_{DF} \in (4.8 \times 10^4 R_0, 2.4 \times 10^3 R_0) \) for both spins, as shown in figures 4(a) and 3(a), respectively. The total resistance, as a parallel connection of two huge spin resistances,
becomes rather high \((2.3 \times 10^3R_D)\). Besides the electric-field induced full-polarization reversal, the electric-field induced n-p junction also contributes to the high resistance state. This is because the transmissions become selective \([37, 38]\). In figure 5(a), we find that the high resistance decreases with an increasing temperature, which implies that the device now supports an insulator state.

The resistances of the same device but with normal ferromagnets \((E_F = 15)\) are plotted in figure 3(b) as a function of \(E_{DF}\). A similar metal state at zero voltage and insulator state at a proper voltage are found. The resulting on-off ratio reads \(1.7 \times 10^4\). However, it is noticed that the \(E_{DF}\) range for the insulator state is different from that in figure 3(a) and it becomes much narrower. In the orange window, it is found that \(E_F - E_{DF} \in (-E_G/2, +E_G/2)\) \((E_{DF} \approx 15.05)\), which means that the Fermi energy enters into the Dirac gap of the free ferromagnet. As a result, electrons of both spins from the device source encounter potential barriers in the free ferromagnet, see figure 1(c2). This is clearly reflected by the rather small transmissions \((< 10^{-4})\) and the rather high resistances \((381R_D\) and \(105R_D)\) for both spins, see figures 4(b) and 3(b), respectively. The case is rather different from the zero voltage case, for which both spins transport over barriers (see figure 1(c1)) and result in near-unit transmissions in figure 4(b) and small resistances \((0.62R_D\) and \(0.57R_D)\) in figure 3(b).

From the above results and discussions, we can see that colossal and extremely large changes in resistance, which require a strong magnetic field in manganite perovskites and other Dirac materials, can be realized by a small gate voltage in the proposed device. The underlying mechanism is that the gate induces either a full-polarization reversal in a half-metallic ferromagnet or an energy gap in a normal ferromagnet, which results in a metal-insulator state of the device. For the former case, the insulator state arises because different spins are blocked in different ferromagnets, while for the latter case, the insulator state arises because both spins are blocked in the free ferromagnet. The energy gap \((1\) meV\) is rather narrow and will be measured out at high temperatures. In the following, we will focus on metal-insulator states stemming from the full-polarization reversal.

Can this remarkable effect be realized in other systems? In figures 3(c)–(f), we plot the calculation results for devices with ferromagnets induced by EuO and EuS. The half-metallic ferromagnet cases based on EuO \((E_F = 7)\) and EuS \((E_F = 9.5)\) are shown in figures 3(c) and (e), and the normal ferromagnet cases based on them \((E_F = 15)\) are shown in figures 3(d) and (f), respectively. It is seen that the device resistance shows similar dependence on the Dirac point of the free ferromagnet as the YIG cases. The resulting CMR are \(3.5 \times 10^4\), \(6.53 \times 10^4\), \(4.1 \times 10^5\), and \(3.18 \times 10^6\) for figures 3(c)–(f), respectively. Meanwhile, the low and high resistances at zero and proper \(E_{DF}\) show similar temperature dependences as the YIG cases (see figure 5(a)). The above behaviors mean that the voltage-induced metal-insulator states and huge changes in resistance can also be realized in devices based on other magnetic insulators.

### 4.2. Ferromagnet-length dependence of the on-off ratios

In all of the above calculations, the ferromagnet length is fixed as \(l_p\); what will happen when it is changed? In figure 6(a), we plot the on-off ratio of the YIG-based device as a function of the lengths of the pinned and free ferromagnets. Surprisingly, it is found that the on-off ratio is rather sensitive to the lengths; it shows a near-exponential dependence on the ferromagnet lengths excepting \(1 < l_D < 1.5\). When the ferromagnet lengths increase from 1 to 2, an extremely large on-off ratio up to \(3.72 \times 10^6\) arises.

The strong enhancement can be understood as the following. The insulator state of the device stems from evanescent transports of both spins, for which the longitudinal wave \(e^{ikx}\) in equation (3) becomes an evanescent wave \(e^{-\kappa x}\), where \(\kappa^2 = -k_m^2\). As a result, the transmission and conductance decreases near-exponentially with an increasing barrier length. Hence, the resistance increases near-exponentially with an increasing length, see the red solid curve in figure 6(b). In contrast, the metal state stems from a ballistic transport of spin
up, for which $e^{k_z}$ is a plane wave. Accordingly, the transmission, conductance and resistance show a weak oscillating dependence on the ferromagnet length, see the red dashed curve in figure 6(b). The enhancement of the on-off ratio follows that of the insulator state.

The cases for devices using EuO and EuS are also plotted in figure 6. It is seen that both the metal-insulator states and the on-off ratio show similar length dependence as the device using YIG. However, the slopes of the profiles are different. For the device using YIG, the change in resistance increases from $3.92 \times 10^5\%$ at $l_{d,f} = 1$ to $3.72 \times 10^8\%$ at $l_{d,f} = 2$; for the device using EuO, the on-off ratio changes from $3.50 \times 10^2\%$ to $5.00 \times 10^6\%$, while for the device using EuS, the values read $4.08 \times 10^5\%$ and $7.14 \times 10^9\%$. An enhancement factor, $\Delta(\log RC)/\Delta l_{d,f}$, can be calculated as $\sim 3, 2$ and 4, respectively. Interestingly, they are found to be proportional to the squares of the smaller effective spin Dirac gaps (but not the squares of the smaller effective spin Dirac gaps), i.e. $\Delta RC(f) \propto e^{-\frac{(\Delta_\uparrow f/\kappa_\uparrow)^2}{\Delta f}}$. This is because the resistance is dominated by the spin with a smaller $\kappa$ (i.e. spin down for all the three graphene ferromagnets, see table 1), and $\kappa^2 \propto -\Delta_\downarrow^2/\epsilon_\uparrow^2$ near the spin Dirac points.

4.3. Temperature dependence of the changes in resistance

Until now, we only considered the zero-temperature on-off ratios. How will it change at high temperatures, especially at the liquid helium and liquid nitrogen temperatures? In figure 5(b), we plot the numerical results for $l = 2$ (the solid red curve) as a function of temperature. As can be seen, the on-off ratio shows a decreasing dependence on the temperature, which is similar to the temperature dependence found in manganite perovskites and other Dirac materials. However, the on-off ratio maintains $305 \times 10^6\%$ (18% smaller than the value at zero temperature) at the liquid helium temperature. This value is hundreds of times higher than the XMR previously reported in other Dirac materials at several K and under magnetic fields of several T [4]. The Curie temperature of the YIG-induced graphene ferromagnet is higher than the room temperature. We have also calculated the on-off ratio at the room temperature by ignoring the inelastic scattering. A change in resistance of 270% is found. The temperature dependence for a device of $l = 1$ is also shown, see the red dashed curve in figure 5(b). It is found that the on-off ratios are much smaller and the temperature dependence is much gentler. Moreover, the higher the temperature, the smaller the difference of the on-off ratios for different ferromagnet lengths. It is noted that the magnetic-field-induced CMR in manganite perovskites exists only near the zero-field transition temperature; the above results show that the proposed electric-field-induced extremely large on-off ratios can survive for a wide range of high temperatures.

The decreasing temperature behavior of the on-off ratio stems from the increased behavior of the low-resistance state and the decreased behavior of the high-resistance state, which is also an important evidence for a metal or insulator behavior. In the following, we will show that these behaviors stem from a negative and positive energy dependence of the zero-temperature conductance around the Fermi energy, respectively, see figure 7. Due to equation (5), a spin-dependent current at a finite temperature $T$ is determined by the zero-temperature currents in an estimated energy range $\sim (E_F - 5T, E_F + 5T)$. In figure 7, we plot the zero-temperature metal-insulator conductances ($G_M$) as a function of energy around the Fermi energy. As can be seen, $G_M$ reaches almost the maximum (minimum) around the Fermi energy, which are direct results of the electronic structure as shown in figure 2. Accordingly, the higher the temperature, the wider the energy range, the smaller (bigger) the finite-temperature $G_M$ and $R_M$. An abnormal decrease of $R_M$ above $T > 58\,K$ is also observed. It stems from a ‘W’-shaped $G_M - E$ profile.

The devices using EuO and EuS show similar temperature dependences (figure 5(b)) and underlying mechanisms (figure 7). However, the temperature ranges are much smaller. At the liquid helium temperature, extremely large changes in resistance up to $4.9 \times 10^6\%$ and $6.5 \times 10^9\%$ are obtained, respectively.
4.4. Metal contacting effects

At last, we consider the influence of metal contacts on the changes in resistance. Several familiar metals, Ag, Cu, Au and Pt at their equilibrium distances with graphene are considered. The contact dopings \((U/E_0)\) equal \(-32, -17, 19\) and \(32\), respectively [39]. The calculated results are shown in figure 8. It is observed that, for devices using YIG and EuS, both the \(R_t\) and \(R_M\) increase as the contacts become non-ideal; the heavier the contact doping, the larger the resistances increase. This is because the contacting resistances are in series with the original ones. However, the low and high resistances increase differently. As a result, the on-off ratio can show different behaviors. It decreases for the devices using YIG (EuS) with positive (negative) dopings, while increases for devices using YIG (EuS) with negative (positive) dopings. Considering that the contact resistance is usually harmful for device performance [40], the latter is a rather interesting and useful result. For devices using EuO, the change in resistance always decreases, slightly for negative contact dopings and sharply for positive contact dopings.

5. Conclusion

In summary, we have proposed a device that can generate extremely large changes in resistance at high temperatures and in a zero magnetic field. The device is composed of a graphene under two YIG strips, where gate voltages are applied. Based on conductance calculations, we have demonstrated that, by applying a proper gate on the free ferromagnet, an on-off ratio up to \(305 \times 10^6\%\) can be obtained at the liquid nitrogen temperature and in a zero magnetic field. This value is hundreds of times higher than the XMR previously observed in Dirac materials at similar temperatures and under magnetic fields of several T. The change in resistance maintains as \(16 \times 10^6\%\) at the liquid nitrogen temperature and in a zero magnetic field, which is still comparable with the CMR observed in manganite perovskites at the same temperature and under magnetic fields of several T. We have indicated that the underlying mechanism for such a remarkable effect is that an electric field induces a reversal of the full polarization in the half-metallic free ferromagnet, which results in a metal-insulator state in the device.

Interesting results also include: (1) the change in resistance shows a near-exponential dependence on the ferromagnet lengths, (2) the longer the two ferromagnets, the sharper the negative temperature-dependence of the on-off ratio, and (3) the effective spin Dirac gap instead of the spin Dirac gap itself plays an important role in the insulator state. We have also shown that the proposed effect can be realized in devices using other magnetic insulators, such as EuO and EuS. Our work should be helpful for developing a realistic switching device. Using an electric field instead of a magnetic field, the proposed device is also far more energy saving and compatible with the ubiquitous voltage-controlled semiconductor technology [41–43].

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