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DOI

[10.1103/PhysRevApplied.16.014032](https://doi.org/10.1103/PhysRevApplied.16.014032)

Publication date

2021

Document Version

Final published version

Published in

Physical Review Applied

Citation (APA)

Zhai, X., Xu, Z., Cui, Q., Zhu, Y., Yang, H., & Blanter, Y. M. (2021). Electrically Controllable Van der Waals Antiferromagnetic Spin Valve. *Physical Review Applied*, 16(1), [014032].
<https://doi.org/10.1103/PhysRevApplied.16.014032>

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Electrically Controllable Van Der Waals Antiferromagnetic Spin Valve

Xuechao Zhai^{1,2,*}, Ziming Xu,³ Qirui Cui,⁴ Yingmei Zhu,⁴ Hongxin Yang^{1b,4,5,†} and Yaroslav M. Blanter^{1b,2,‡}

¹*Department of Applied Physics, Nanjing University of Science and Technology, Nanjing 210094, People's Republic of China*

²*Kavli Institute of NanoScience, Delft University of Technology, Delft 2628 CJ, Netherlands*

³*School of Science, Nanjing University of Posts and Telecommunications (NJUPT), Nanjing 210023, People's Republic of China*

⁴*Ningbo Institute of Materials Technology and Engineering, Chinese Academy of Sciences, Ningbo 315201, People's Republic of China*

⁵*Center of Materials Science and Optoelectronics Engineering, University of Chinese Academy of Sciences, Beijing 100049, People's Republic of China*

(Received 28 February 2021; revised 22 June 2021; accepted 24 June 2021; published 13 July 2021)

We propose a spin valve that is based on van der Waals antiferromagnetism and is fully electrically controlled. The device is composed of two antiferromagnetic terminals that allow for vertical bias control and a linked central scattering potential region. The magnetoresistance varies significantly when the bias orientations in two terminals are switched from parallel to antiparallel, because this switch induces a mismatch of the bands for the same spin projection in different parts of the system. It is shown from density-functional calculations that bilayer graphene encapsulated by two atomic layers of $\text{Cr}_2\text{Ge}_2\text{Te}_6$ provides a material platform to realize the antiferromagnetism, which is robust against the required vertical electric fields.

DOI: [10.1103/PhysRevApplied.16.014032](https://doi.org/10.1103/PhysRevApplied.16.014032)

I. INTRODUCTION

The spin valve, consisting of two or more layered conducting ferromagnetic (FM) materials, is a key element of spintronics [1,2]. The electrical resistance of a spin valve switches between high and low values depending on whether the alignment of the magnetization in the layers is parallel (P) or antiparallel (AP). The mechanism of this change is the giant magnetoresistance (MR) effect [3], which is the basis of modern commercial devices such as reading heads for computer hard drives or magnetic sensors. In the past decade, the spin valve has also attracted much attention in two-dimensional (2D) materials (see, e.g., Refs. [4–10]), which have easily tunable electronic properties.

Recently discovered 2D van der Waals (vdW) magnets, notably $\text{Cr}_2\text{Ge}_2\text{Te}_6$ (CGT) [11] and CrI_3 [12], have easily manipulated magneto-electric and magneto-optical properties [13,14] and therefore are very promising for spintronics. In particular, the layered antiferromagnetic (LAF, i.e., two adjacent FM monolayers coupled

antiferromagnetically) order confirmed in bilayer insulating CrI_3 [12,15–17] adds an intriguing dimension to the development of 2D antiferromagnetic spintronics with low power consumption and robustness against perturbations from magnetic fields [18,19]. A few proposals exploiting LAF have been put forward, including the manipulation of electron spin by a high-intensity electric field [20] and exotic topological phases by the electro-optical effect [21]. However, the intrinsic LAF systems usually have large band gaps, no or weak spin polarization, or low electron mobility [12–17,20], which strongly limits their applications in the design of electronic circuits with low energy consumption.

Recent progress has made it possible to fabricate 2D structures composed of almost any layers on top of each other [22–24]. These structures display a higher variety of properties than natural materials, can be tailored to have predefined physical properties, and open a way to the design of high-performance devices based on unusual interface physics. Motivated by the latest advances in the fabrication of four-layer vdW heterosystems [typically graphene/bilayer CrI_3 /graphene [8] and WSe_2 /bilayer graphene (BLG)/ WSe_2 [25]] and a BLG-based device that allows for local bias control [26], we propose an artificial four-layer LAF device [Fig. 1(a)] with

*zxcnust@163.com

†hongxin.yang@nimte.ac.cn

‡Y.M.Blanter@tudelft.nl

FIG. 1. (a) Top: a sketch of a LAF spin valve based on BLG encapsulated by 2D FM insulators. 1 (2) marks the layer of BLG. The device consists of two LAF terminals [left (L), right (R)] and a linked central region (C, length d), where a collinear DW (the magnetic periodicity along y is not shown) is taken as an example. Local dual gates in L and R are used to induce the vertical electric field E_z and tune the Fermi energy E_F [8,26]. Bottom: typical bands are plotted to show how a high MR is achieved for the P arrangement of E_z in L and R. (b) A general picture of a field-modulated layer- and the spin-resolved density of states (DOS) for the case on the right. By shifting the spin-layer-dependent states in energy, the LAF field in the model given in Eq. (1) induces a band gap in the BLG and then E_z fully drives the spin polarization within the gap.

two gated terminals linked by the central scattering region to realize a fully electrically controlled spin valve. Its operation mechanism is similar to that of existing giant MR devices [3–10], with the difference that it does not require a magnetic field: by altering the mutual bias orientation in the two terminals, from P to AP, the MR of in-plane transport can be varied significantly. Due to the magneto-electric coupling [Fig. 1(b)], the spin-valve effect is transferred to the nonmagnetic vdW layers. We start with a phenomenological model and further use density-functional theory (DFT) calculations to demonstrate the concept with a material platform—BLG encapsulated by two single layers of CGT.

Our proposed LAF spin valve has three advantages over the usual spin valves [3–10]: (i) It is robust against environmental perturbations from magnetic or null-stray fields due to antiferromagnetism [18–20]; (ii) it requires only a small vertical bias to control, is reversible to operate, and is superior to traditional magnetic control [3]; and (iii) it is

not sensitive to the presence of domain walls (DWs), thus enabling the device to possess strong antijamming of DWs. The two latter points are not trivial and we demonstrate their validity here.

II. MODEL AND THEORY

Our model of four vdW layers in Fig. 1(a) contains the nonmagnetic conducting BLG with Bernal stacking [27,28] in the middle and two FM insulating layers located, respectively, on the upper and lower sides. As a typical conducting system, BLG has the advantages of highly adjustable electron mobility [27,29] and a long spin relaxation length [30–32], which can be very little affected by the proximity of vdW [25,33] and suggest an easily prepared platform to realize low-consumption devices. We assume that the two insulating layers are identical 2D magnets with out-of-plane magnetization (favorable for ultrahigh-density data storage [34]) and do not contribute to the low-energy electronic bands except for proximity-induced magnetism. An effective low-energy Hamiltonian from an empirical tight-binding model (see Appendix A) reads

$$H_0 = v_0(p_x\sigma_x + \xi p_y\sigma_y) - \frac{\gamma}{2} \sum_{i=x,y} \tau_i\sigma_i + m_z\tau_zs_z + \lambda_U\tau_z, \quad (1)$$

where σ_i , τ_i , and s_i ($i = x, y, z$) are, respectively, the Pauli matrices for sublattice isospin, layer isospin, and real spin. Here, $\xi = \pm 1$ indicate two opposite valleys, K and K' , of the hexagonal Brillouin zone, $p_{x(y)}$ is the momentum component, $v_0 = 3at/2$ is the Fermi velocity in graphene (a and t are the bond length and nearest-neighbor hopping, respectively), and γ , m_z , and λ_U denote the strengths of the interactions. The first and second terms are, respectively, the massless Dirac term and the interlayer nearest-neighbor coupling (usually sufficient to capture the band features [27,28]). The third term indicates the LAF field induced by short-range magnetic proximity from two FM layers [9,35]. The last term represents the vertical bias.

We also consider other possible interactions that might be present but that do not qualitatively change the MR effect with which we are concerned. They read [25,33,36] $H' = H'_\Delta + H'_R$, where $H'_\Delta = \Delta\sigma_z$ represents the intralayer inversion-symmetry breaking induced by interface lattice mismatch and $H'_R = \lambda_R(\sigma_xs_y - \xi\sigma_ys_x)/2$ [36–38] is the proximity-induced Rashba interaction. For simplicity, we take H' to be the same in both monolayers. Disregarding H'_R , the four low-energy bands are solved as follows:

$$E^2 = v_0^2p^2 + \frac{\gamma^2}{2} + \alpha^2 + \Delta^2 - \frac{1}{2}\sqrt{4v_0^2p^2(4\alpha^2 + \gamma^2) + (\gamma^2 - 4\alpha\Delta)^2}, \quad (2)$$

where $\alpha = \lambda_U + sm_z$ ($s = \pm 1$ denote spin up and spin down, respectively). In the actual numerical calculations below (Fig. 2), weak H'_R is also taken into account.

As sketched in Fig. 1(a), the system is divided into three regions, including the left (L) and right (R) terminals and the central (C) scattering region. We use an effective empirical potential [39] $\lambda_U(x) = U \text{erf}(x/d)$ and U , respectively, to simulate the AP and P bias configurations. Here, “erf” is the error function and $x = 0$ is the center of region C.

Depending on the sample quality, the system can be a pristine LAF crystal when no DW is present in region C ($dm_z/dx = 0$). The simplest DW is collinear, as sketched in Fig. 1(a), and can be simulated by a smooth potential [40]

$$m_z(x) = M_0 \tanh \frac{\pi x}{d}. \quad (3)$$

Here, the magnetization m_z vanishes at $x = 0$ due to the destructive overlap between opposite spins. We can further consider the noncollinear Bloch ($m_y^2 = M_0^2 - m_z^2$, $m_x = 0$) and Néel ($m_x^2 = M_0^2 - m_z^2$, $m_y = 0$) types of DW due to spin fluctuation [40] by adding in-plane magnetization terms $m_y \tau_z s_y$ and $m_x \tau_z s_x$ to Eq. (1), respectively. The in-plane magnetization of Bloch (Néel) type is parallel (perpendicular) to the DW, and the sign of m_x or m_y indicates the DW chirality. The spin valve functions qualitatively in the same way with or without DWs; however, the details of the scattering potential differ.

Taking the collinear DW in Fig. 1(a) for example, we illustrate typical band-to-band tunneling to realize the high-MR state for the P bias configuration due to spin mismatch between L and R. By tuning the direction of the vertical bias, the LAF spin valve can be switched on or off. The giant MR effect is understood from a general picture of phase transition in Fig. 1(b), seen from the model given in Eq. (1), in which the LAF field drives the BLG to be insulating by moving the states with opposite spins in

the same layer and the states with the same spin in opposite layers in opposite directions. The vertical bias induces full spin polarization by moving the states in the opposite layers in opposite directions. More crucial for the realization of the entirely electric control of the spin-valve effect is that the opposite bias induces the opposite spin polarization. We characterize the MR effect by a resistance ratio [3,7],

$$\rho_{\text{MR}} = \frac{R_{\text{AP}} - R_{\text{P}}}{R_{\text{AP}} + R_{\text{P}}}, \quad (4)$$

where $R_{\text{P(AP)}} = 1/G_{\text{P(AP)}}$ is the resistance for the P (AP) case. The conductance can be calculated by using the nonequilibrium Green’s function method [41] as

$$G = \frac{e^2}{h} \int_{-\infty}^{+\infty} dE \mathcal{T}(E) \left(-\frac{\partial}{\partial E} f(E - E_F) \right), \quad (5)$$

with the transmission given by $\mathcal{T}(E) = \text{Tr}(\mathcal{G} \Gamma_L \mathcal{G}^\dagger \Gamma_R)$, where $\mathcal{G}$ is the Green’s function of the system, “Tr” denotes the trace, $\Gamma_{L,R} = i(\Sigma_{L,R} - \Sigma_{L,R}^\dagger)$ represents the broadening of L and R, and $f(E - E_F)$ is the Fermi-Dirac distribution function.

III. RESULTS AND DISCUSSION

We first solve the model by performing calculations in an armchair-edged sample with width $W = 40\sqrt{3}a$. The two terminals, L and R, are assumed to be semi-infinite. A relatively short scattering region with length $d = 10a$ is fixed and the main electrically tunable parameters are the bias voltage U and the Fermi energy E_F (see details in Appendix B and Ref. [42]). We keep $|U|$ the same for the P and AP configurations and fix the parameters $t = 2.7$ eV and $\gamma = 0.39$ eV from bare GBL [28] and $M_0 = 27$ meV, $\Delta = 13.5$ meV, and $\lambda_R = 0.9$ meV.

FIG. 2. The calculated results from the lattice model (see Appendix A) with parameters $t = 2.7$ eV, $\gamma = 0.39$ eV [28], $M_0 = 27$ meV, $\Delta = 13.5$ meV, and $\lambda_R = 0.9$ meV. (a) The MR ratio versus E_F for different DW types (no DW, collinear, Bloch, and Néel) for an armchair-edged strip with $W = 40\sqrt{3}a$ under $U = -M_0$ and $T = 0$. The symbols I–IV label four MR states with clear boundaries. (b) The contour plot of the MR ratio in the (E_F, U) plane for the armchair-edged strip with collinear DW at $T = 0$. The dashed line corresponds to the cross section $U/M_0 = -1$ used in (a). (c) The band gap E_g in terminal L (green dashed line). For comparison, the spin-dependent band gap E_g^s is shown under $\lambda_R = 0$ (red circle for spin up, black rectangle for spin down). The blue circles denote four gap values $E_{g,i}$ ($i = 1$ to 4) under $U/M_0 = \pm 1$. (d) The influence of the temperature on the MR ratio versus E_F under $U = -M_0$ for the collinear-DW case in (a).

When $U = -M_0 = -27$ meV is fixed ($E_z \simeq -0.5$ V/nm [29]), our results for the MR ratio versus E_F in Fig. 2(a) show that a nearly perfect MR effect, i.e., $|\rho_R| \simeq 100\%$, appears for $|E_F|/M_0 \in (0.5, 1.5)$, independent of whether or not the DW is present. The MR ratio of this spin-valve effect changes sign if a DW appears. This is because the m_z -induced spin polarization near E_F changes sign in terminal R (but is unchanged in L). Nontrivially, we find that the giant MR effect is not sensitive to whether the DW type is collinear, Bloch, or Néel, and meanwhile, the spin-valve effect is irrelevant to the magnetic chirality of DW, which can be confirmed by changing the sign of m_x or m_y in the calculations. To illustrate the MR effect in more detail, we show the contour plot of the MR ratio in the (E_F, U) plane for the collinear DW in Fig. 2(b). The result shows that there is a sharp transition between the absence of the MR effect for lower values of E_F and the nearly perfect MR effect for higher values of E_F [see also $|E_F|/M_0 < 0.5$ in

Fig. 2(a)]. This parameter regime of no MR, $G_{P(AP)} \equiv 0$, is attributed to either of the following two factors: one is the absence of transport states induced by the band gap in the two terminals and the other is the spin mismatch for both the P and the AP bias configurations.

To precisely clarify the result in Fig. 2(b), we plot the band gap E_g (green dashed line) for terminal L as a function of U in Fig. 2(c). For contrast, we also plot the spin-dependent band gap E_g^s (circle for spin up, rectangle for spin down) by ignoring λ_R in the calculations. It is found that the total band gap changes very little if λ_R is added due to $\lambda_R \ll M_0, \Delta$. In the presence of the DW in Eq. (3), the band gap in terminal R can be directly obtained by interchanging the tabs for opposite spins in Fig. 2(c) due to the opposite magnetic orientation in the two terminals. Thus, it is sufficient to obtain all the information for the band gap by analyzing either of the two terminals. After sorting four values $E_g^{s=\pm}(\pm U)$ for

FIG. 3. The DFT results for CGT/BLG/CGT. (a) The side view of the optimized unit cell. (b) The vertical electric field E_z tuned energy difference, on average, per atom between the FM and LAF states. (c) The energy bands for $E_z = 0$ (left) and ± 0.2 V/nm (right). (d) The spin-dependent band gap as a function of E_z .

terminal L in order from lowest to highest, labeled as $E_{g,i}$ ($i = 1, \dots, 4$), we find, considering the tunneling mechanism in Fig. 2(b), that MR vanishes (I, $\rho_{\text{MR}} = 0$) for $|E_F| < E_{g,2}/2$, and that the nearly perfect MR (II, $|\rho_{\text{MR}}| \geq 95\%$) occurs under the conditions $|E_F| \in (E_{g,2}, E_{g,3})/2$ for $U < 0$ and $|E_F| \in (E_{g,2}, E_{g,4})/2$ for $U > 0$. Likewise, the normal giant MR (III, $95\% > |\rho_{\text{MR}}| \geq 25\%$ [3]) holds for $|E_F| \in (E_{g,3}, E_{g,4})/2$ and the weak MR (IV, $0 < |\rho_{\text{MR}}| < 25\%$) happens for $|E_F| > E_{g,4}/2$.

Without a DW, the dependence of $E_g(U)$ is identical for the two terminals and the condition for the MR ratio to vanish in Fig. 2(a) is strictly given by $|E_F| < E_g(U)/2$ [see the blue region in Fig. 2(c)]. Independent of whether or not the DW is present, the nearly perfect MR can work for the overlapping band-gap region of opposite spins [see the yellow region in Fig. 2(c)] by tuning the magnitude of the bias, which can be different for L and R, in addition to the sign of the bias.

Figure 2(d) further considers the influence of the temperature, showing the MR ratio at $T = 60$ K and 120 K, for the collinear DW at $U = -M_0$. Unsurprisingly, the plot of the MR ratio versus E_F becomes smooth due to the presence of the Fermi-Dirac distribution in Eq. (5), in comparison with that in Fig. 2(a). The energy broadening induced by the temperature can cover the whole band-gap region and thus enables the giant MR effect when E_F lies inside the band gap. Note that the pattern of the curves in Fig. 2(d) remains the same even if the T -modified M_0 is considered in a specific system. This means that the proposed spin-valve mechanism remains operational up to the Curie temperature of the FM layers.

In particular, the numerical result E_g^s for the finite-size strip in Fig. 2(c) agrees with the analytical result for the 2D bulk case $E_g^s(U) \simeq 2(\gamma^2/2 + \alpha^2 + \Delta^2 - |\gamma^2 - 4\alpha\Delta|/2)^{1/2} = 2|U + sm_z + \Delta|$ derived from Eq. (2) for $\gamma^2 > 4\alpha\Delta$. This reveals a negligible edge effect. Further calculations indicate that the edge effect on band gap can be ignored when W exceeds a critical value $20\sqrt{3}a$ (approximately 5 nm), consistent with the experimental report [43] for bare BLG. Additionally, it is argued that the electron mean free path can be estimated by $\ell \simeq \lambda_F = \hbar v_0 / \sqrt{|E|\gamma}$ [39] (where λ_F is the Fermi wavelength and $\ell \simeq 248a = 35.2$ nm for $E = 0.01t$). The spin diffusion length may be estimated by $\ell_s^{-1} \simeq 2\ell[\lambda_R/(\hbar v_F)]^2$ (the Fermi velocity $v_F = 2v_0\sqrt{|E|/\gamma}$) in the Dyakonov-Perel mechanism [44]. The ballistic spin transport, on which we deliberately focus here, should always dominate for $d \ll \ell, \ell_s$ (see Appendix C; ℓ_s is usually on the micrometer scale). The nanoscale W and d here should be experimentally available [40,43].

For zigzag-edged systems, the LAF field or vertical bias can induce edge states [21], which reduce the MR ratio. However, an edge disorder or defects can increase the MR ratio by hindering the edge transport. For both the armchair or zigzag-edged systems with $W < 20\sqrt{3}a$, the spin-valve

effect can still hold in spite of strong quantum confinement, which only leads to some quantitative variations in MR.

IV. PROPOSAL FOR MATERIALS REALIZATION

We now propose, based on DFT calculations [45–49], that BLG encapsulated by CGT monolayers, each of which is a FM semiconductor [11,14], is feasible for the realization of antiferromagnetism in our model. There has been a previous LAF prediction in $\text{CrI}_3/\text{BLG}/\text{CrI}_3$ [9] but the formation of mixed low-energy bands from both BLG and CrI_3 strongly limits its superiority. For the compounds of graphene and CGT, the influence of the spin-orbit couplings on the low-energy bands is much smaller than the other parameters and can be disregarded, as demonstrated previously [50,51].

In Fig. 3, we show the DFT results for CGT/BLG/CGT with an average equilibrium interlayer spacing $d = 3.65$ Å between the CGT monolayer and its close-contacted graphene monolayer. The unit cell in Fig. 3(a) is constructed with a 5×5 supercell of 50 carbon atoms per graphene layer and a $\sqrt{3} \times \sqrt{3}$ supercell of 12 chromium atoms per CGT layer. (for other computable details, see Appendix D). To clarify the magnetism of the ground state, we show the energy difference, on average, per atom (ΔE) between the FM and LAF states in Fig. 3(b), which strongly suggests that the system has a LAF ground state that is robust against the small E_z on which we focus. In contrast, the bilayer CGT itself is FM [11], which becomes LAF here when the layers are spatially separated.

We show the DFT band at $E_z = 0$ in the left panel of Fig. 3(c), where the inset shows an enlargement of the local band near E_F and is well fitted by $\Delta = 8.9$ meV, $m_z = -7.2$ meV, $t = 2.441$ eV, and $\gamma = 0.396$ eV in our model and an additional interlayer next-nearest-neighbor hopping $\gamma' = 0.267$ eV. Here, we have $m_z < 0$ and $\Delta > |m_z|$, different from the assumptions in Fig. 2. As expected, we show the contrasting spin polarization of the bands for $E_z = \pm 0.2$ V/nm, which is much smaller than the $1 \sim 3$ V/nm experimentally accessible in vdW systems of BLG [29] and bilayer CrI_3 [16]. We plot the spin-dependent band gap E_g^s versus E_z in Fig. 3(d), where the spin-up and spin-down band gaps vanish for $E_z > 0.23$ V/nm and $E_z < -0.1$ V/nm, respectively. This is beyond our model, because CGT makes a non-negligible contribution to E_g^s as E_z increases.

V. CONCLUSION

Our proposal for an antiferromagnetic spin valve based on the vdW magnetic proximity effect paves the way for engineering spintronic devices with low energy consumption by making full use of magneto-electric coupling. The appearance of nearly perfect MR and the robustness of the antiferromagnetism against the required vertical bias reveal the superiority of our proposed device.

Other adjustable quantities, e.g., the stacking order, inter-layer sliding, rotation angle or layer number, may bring rich opportunities for developing 2D antiferromagnetic spintronics.

ACKNOWLEDGMENTS

This work was supported by the National Natural Science Foundation of China (NSFC) with Grants No. 12074193, No. 61874057, and No. 11874059, the QingLan Project of Jiangsu Province (2019), the Key Research Program of Frontier Sciences, Chinese Academy of Sciences (CAS), Grant No. ZDBS-LY-7021, the Beijing National Laboratory for Condensed-Matter Physics, and the Natural Science Foundation of Zhejiang Province (Grant No. LR19A040002). We are grateful to S. Jiang for insightful experimental discussions.

APPENDIX A: LATTICE HAMILTONIAN

According to Refs. [27,28], for BLG, the low-energy effective Hamiltonian in Eq. (1) can be empirically derived from the lattice Hamiltonian

$$H_0 = -t \sum_{\langle ij \rangle_{\parallel} \alpha} c_{i\alpha}^{\dagger} c_{j\alpha} - \gamma \sum_{\langle ij \rangle_{\perp} \alpha} c_{i\alpha}^{\dagger} c_{j\alpha} + m_z \sum_{i\alpha} \mu_i c_{i\alpha}^{\dagger} s_z c_{i\alpha} + \lambda_U \sum_{i\alpha} \mu_i c_{i\alpha}^{\dagger} c_{i\alpha}, \quad (\text{A1})$$

where $\alpha(\beta)$ denotes spin up (down), $\mu_i = \pm 1$ holds for the i site on the top (bottom) layer, s_z is the z component of the Pauli matrix, $c_{i\alpha}^{\dagger}$ creates an electron with spin polarization α at site i , $\langle i, j \rangle$ runs over all the nearest-neighbor-hopping (NNH) sites, and the subscript $\parallel$ ($\perp$) means in-plane (out-of-plane). The interlayer coupling between the next NNH sites can be considered by adding $-\gamma' \sum_{\langle \langle ij \rangle \rangle_{\perp} \alpha} c_{i\alpha}^{\dagger} c_{j\alpha}$ [28], which is needed for the quantitative comparison of the DFT bands and the model. Usually, other weaker hopping elements between atoms at greater distances are not essential for the problem at hand and are therefore disregarded.

With regard to the vdW proximity effect between two different materials, we need to further consider the following perturbation H' in lattice space:

$$H' = \Delta \sum_{i\alpha} \xi_i c_{i\alpha}^{\dagger} c_{i\alpha} - i\lambda_R \sum_{\langle ij \rangle_{\parallel} \alpha\beta} c_{i\alpha}^{\dagger} (\mathbf{s} \times \mathbf{d}_{ij})_{\alpha\beta}^z c_{j\beta}, \quad (\text{A2})$$

where the first term represents the intralayer inversion-symmetry breaking induced by the interface-lattice mismatch [$\mu_i = \pm 1$ holds for sublattice A (B) in each monolayer] and the second term represents the proximity-induced Rashba spin-orbit coupling [25,33,36]. As is allowed by symmetry, no difference between the two layers of BLG for each term of H' is considered.

APPENDIX B: DETAILS OF ELECTRICAL CONTROL

The local-bias controllability of the four-layer vdW geometry can be analyzed by a capacitance model, as shown in Fig. 4. For simplicity, here we ignore the vacuum layer and the additional nonmagnetic (such as hexagonal boron nitride) dielectric layers usually used in a realistic device [8,15,16,26]. The top and bottom layers of BLG experience an electric field E_t and E_b , respectively, given by $E_t = -V_t/D_2$ and $E_b = V_b/D_2$. The contribution of E_t and E_b to the electric field in the BLG can be estimated by [42] $E_{av} = (E_t + E_b)/2 = (-V_t + V_b)/(2D_2)$. The total effective electric field in the BLG reads $E_z = E_{av} - (\rho_1 - \rho_2)\mathcal{S}/(2\epsilon_1)$, where $\mathcal{S}$ denotes the effective area of the capacitor plate. Here, ρ_1 and ρ_2 denote the charge densities (with respect to the neutral case) accumulated on layers 1 and 2 of the BLG, respectively; $\rho_1 - \rho_2 \equiv \Delta\rho$ is the difference in the surface charge densities between two layers of BLG, and can be obtained by self-consistent DFT calculations [42] after fixing V_t and V_b . The directly observable electrical parameters V_t and V_b can be equivalently translated into the effective interlayer bias voltage $2U$ [for layer 1 (2), the potential energy is $\pm U$] and the Fermi energy E_F by

$$2U = -D_1 e E_z, \quad E_F = (\rho_1 + \rho_2)\mathcal{S}_0/(2N_0), \quad (\text{B1})$$

where e is the electron charge ($|e| = -e$) and $\mathcal{S}_0$ and N_0 are the area and carbon-atom number of a unit cell, respectively.

FIG. 4. A simple capacitance model for the four-layer vdW geometry of a dual-gated electrode [this applies equally to the two electrodes in Fig. 1(a)]. D_1 and D_2 label the vertical distances. ϵ_1 and ϵ_2 denote the out-of-plane dielectric constants for the BLG and the ferromagnetic (FM) insulator, respectively. ρ_1 and ρ_2 denote the charge densities (with respect to the neutral case) accumulated on layers 1 and 2 of the BLG, respectively. $V_{t(b)}$ represents the effective voltage added to the interface between the top (bottom) FM layer and the neighboring gate.

APPENDIX C: INFLUENCE OF STRUCTURE PARAMETER d ON MR

In Fig. 5, we show the influence of the length d of the central region on the MR ratio for the collinear-DW case in the device of Fig. 1(a), where the P and AP bias configurations correspond to the high-resistance and low-resistance states, respectively. To understand the role of d in electron tunneling, we need to use two length parameters, which are the electron mean free path ℓ and the spin diffusion length ℓ_s .

Two specific cases are considered, including (i) $(U, E_F) = (-M_0, M_0)$ (open square) and (ii) $(U, E_F) = (M_0, 2M_0)$ (solid circle). Our results show that the MR ratio is essentially unchanged for smaller d (a sharp DW interface), because the ballistic spin transport determined by $d < \ell, \ell_s$ dominates. Here, ℓ can be approximately estimated by [39] $\ell \simeq \lambda_F = \hbar v_0 / \sqrt{|E|\gamma}$, which is about $248a \simeq 35.2$ nm and $175a \simeq 24.9$ nm for cases (i) and (ii), respectively. Note that λ_F is the Fermi wavelength in BLG, E is the electron energy, and $v_0 = 3at/(2\hbar)$ is the Fermi velocity in monolayer graphene (a is the bond length). In the framework of the Dyakonov-Perel mechanism [44], ℓ_s can be estimated by $\ell_s^{-1} \simeq 2\ell[\lambda_R/(\hbar v_F)]^2$, where $v_F = 2v_0\sqrt{|E|\gamma}$ is the Fermi velocity of the low-energy electrons ($|E| \ll \gamma$) in BLG. It is estimated that $\ell_s \simeq 1.65$ μm and 2.49 μm for cases (i) and (ii), respectively. As d increases, the transmission decreases for the low-resistance state (AP bias configuration) due to the enhanced DW scattering and, meanwhile, the transmission increases slightly for the high-resistance state (P bias configuration) due to the increased weak spin-flipping scattering from the Rashba interaction. The reduced value of the MR ratio is no more than 5% because $d/\ell_s < 3 \times 10^{-3}$ is valid for both cases (i) and (ii) under the condition of $d < 300a = 42.6$ nm.

FIG. 5. The dependence of the MR ratio on the length d of the central region for the collinear DW case in the device of Fig. 1(a). The electrical parameters read $(U, E_F) = (-M_0, M_0)$ (solid line) for case (i) and $(M_0, 2M_0)$ (dashed line) for case (ii). The fixed parameters read $t = 2.7$ eV, $\gamma = 0.39$ eV, $M_0 = 27$ meV, $\Delta = 13.5$ meV and $\lambda_R = 0.9$ meV, as used in Fig. 2.

It is strongly argued that ℓ affects the size of the ballistic transport regime [28], while ℓ_s dictates the scale of the spin-flipping scattering. Actually, when d exceeds ℓ (much less than ℓ_s here), the disorder in a real sample cannot be neglected in the strict sense, leading to dissipative transport beyond the ballistic regime on which we focus. Undoubtedly, the ballistic spin transport should always dominate for $d \ll \ell, \ell_s$.

APPENDIX D: DFT CALCULATIONS

Our first-principles calculations are carried out within the framework of DFT as implemented in the Vienna *ab initio* simulation package (VASP) [45–47]. The generalized gradient approximation (GGA) based on the Perdew-Burke-Ernzerhof (PBE) function is used to describe the exchange correlation energy [48]. The $\text{Cr}_2\text{Ge}_2\text{Te}_6/\text{BLG}/\text{Cr}_2\text{Ge}_2\text{Te}_6$ (CGT/BLG/CGT) heterostructure is constructed by depositing a $5 \times 5 \times 1$ supercell of BLG with *AB* stacking on a $\sqrt{3} \times \sqrt{3} \times 1$ CGT supercell, with a small lattice mismatch of 0.38% (according to the same settings of lattice constants as in the case of single-interface BLG/CGT [50]).

We apply the DFT-D3 method in all calculations to treat the vdW interactions [49]. A vacuum region larger than 15 Å is adopted to avoid the interaction between periodic replicas. The cutoff energy to expand the plane wave is set to 520 eV. Γ -centered k meshes of $4 \times 4 \times 1$ and $6 \times 6 \times 1$ are adopted in the calculations of the structural relaxation and the charge distribution, respectively. Considering the strongly correlated Cr-3D electrons, the GGA+U method is adopted ($U = 1$ eV) [11]. All structures are fully relaxed until the Hellmann-Feynman force on each atom is less than 10^{-2} eV/Å and the convergence criterion of the electronic interaction is set to 10^{-6} eV.

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