Layer-Magnetization-Tuned Topological Phases in Mn$_2$Bi$_2$Te$_5$ Films

Yue Li, Yadong Jiang, Jinlong Zhang, Zhaochen Liu, Zhongqin Yang, and Jing Wang

1 State Key Laboratory of Surface Physics and Department of Physics, Fudan University, Shanghai 200433, China
2 Key Laboratory of Computational Physical Sciences (Ministry of Education), Fudan University, Shanghai 200433, China
3 Collaborative Innovation Center of Advanced Microstructures, Nanjing 210093, China
4 Institute for Nanoelectronic Devices and Quantum Computing, Fudan University, Shanghai 200433, China

(Dated: January 20, 2020)

The interplay between band topology and magnetic order could generate a variety of time-reversal-breaking gapped topological phases with exotic topological quantization phenomena, such as quantum anomalous Hall (QAH) insulators and axion insulators (AxI). Here by combining analytic models and first-principles calculations, we predict QAH and AxI phases can be realized in thin film of an intrinsic antiferromagnetic van der Waal material Mn$_2$Bi$_2$Te$_5$. The phase transition between QAH and AxI is tuned by the layer magnetization, which would provide a promising platform for chiral superconducting phases. We further present a simple and unified continuum model that captures the magnetic topological features, is generic for Mn$_2$Bi$_2$Te$_5$ and MnBi$_2$Te$_4$ models.

The discovery of time-reversal-invariant topological insulator brings the opportunity to realize a large family of exotic topological quantization phenomena [1–4]. The interplay between band topology and magnetism could give rise to a variety of exotic time-reversal-breaking gapped topological states, including the quantum anomalous Hall (QAH) effect with dissipationless chiral edge states [5–8], and axion insulator (AxI) displaying topological magnetoelectric effects [8–14], and chiral superconducting state with Majorana fermions (if in proximity to superconductors) [15–17]. Interestingly, the QAH effect may find applications in low-power-consumption electronics and non-Abelian braiding of Majorana fermions is useful in topological computation [18–20]. Despite of these predicted important physical effects, until now only a few of them have been experimentally proved, due to a limited number of magnetic topological insulator (TI) materials. A prime example is the QAH effect experimentally observed in magnetically doped (Bi,Sb)$_2$Te$_3$ film [21–24]. However, the random magnetic dopants limit the quality and exchange gap [25] of the material, which further constrain the quantization of AHE appearing only at very low temperatures. In proximity with an s-wave superconductor, such a strongly disordered QAH system at coercivity by the random magnetic domains complicates the transport experiments in millimeter-size sample [26–31]. Therefore, finding stoichiometric TI with an innate magnetic order are highly desired, which would provide a homogenous platform for high temperature QAH effect and coherent chiral Majorana fermions.

The first intrinsic magnetic TI MnBi$_2$Te$_4$ (MBT) discovered recently is an interesting candidate for observing these topological phenomena [32–38]. For instance, the zero-field QAH effect has been observed at an elevated temperature [36]. Given the importance of magnetic TIs as a platform for new states of quantum matter, it is important to search for new material systems that are stoichiometric crystals with well-defined electronic structures, preferably with simple surface states, and describable by simple theoretical models. In this work, by combining analytic models and first-principles calculations, we predict layer magnetization tuned AxI and QAH in new magnetic TI family Mn$_2$Bi$_2$Te$_5$ (MBT) films. Its bulk hosts a dynamical axion field, which was first proposed in Ref. [39] focusing on the interplay of bulk antiferromagnetic (AFM) fluctuations and axion electrodynamics. Here we will focus on the various topological states in its thin film form with different static magnetic ordering.

MBT is a layered ternary tetradyminite compound that consists of ABC stacking Te$_1$-Bi$_1$-Te$_2$-Mn$_1$-Te$_3$-Mn$_1$'-Te$_2'$-Bi$_1'$-Te$_1$' nonuple layers (NL), which has been successfully synthesized in experiments recently [40]. It has a hexagonal crystal structure shown in Fig. 1(a) with space group $P\overline{3}m1$ (No. 164), which can be viewed as layered TI Bi$_2$Te$_3$ with each of its Te-Bi-Te-Bi-Te quintuple layer intercalated by two additional Mn-Te bilayers. The trigonal axis (threefold rotation symmetry $C_{3z}$) is defined as the z axis, a binary axis (twofold rotation symmetry $C_{2z}$) is defined as the x axis and a biaxial axis (in the reflection plane) is defined as the y axis for the coordinate system. The system has inversion symmetry $P$ with Te$_3$ site as an inversion center if the spin moments of Mn are ignored.

As far as the magnetic order is concerned, it appears that the Mn spins couple ferromagnetically in each layer, but the adjacent Mn layers couple antiparallel to each other. The ferromagnetic (FM) order in each Mn layer can be understood from the Goenough-Kanamori 90$^\circ$ rule, while the AFM coupling between adjacent Mn layers is from the interlayer superexchange similar to MBT. The local magnetic moments are roughly 4.59$\mu_B$ independent of the film thickness. Table I lists the thickness dependence of magnetism, and the magnetic anisotropic energy (MAE) for 2 to 7 NL as well as bulk are positive and insensitive to layer thickness, indicating the Néel-type AFM order along z axis is likely the ground state. The non-collinear and other possible collinear magnetic con-
configuration of the symmetry operations is given by $C_3$ are the three-fold rotation symmetry
superscripts $+^\ast$, $−^\ast$ give the orbital-projected band structures without and with SOC, respectively. (e) Schematic diagram of the band inversion induced by SOC at $\Gamma$. The green dashed line represents the Fermi level. (f) The energy and momentum dependence of the local density of states (LDOS) on the (110) and (001) surface, respectively.

Then we turn to the electronic and topological properties of the material. To have an intuitive understanding of the underline physics, we start with the bulk electronic structure. The detailed methods can be found in the Supplemental Material [42]. For the AFM ground state, the band structures without and with spin-orbit coupling (SOC) are shown in Figs. 1(c) and 1(d), respectively. Mn $d$-bands are far away from the band gap due to a large spin splitting (>$5\,eV$), and only Bi/Te $p_z$-bands are close to the Fermi level with an anticrossing feature around the $\Gamma$ point from the band inversion, suggesting the nontrivial topology in bulk M$_2$BT. To characterize the low-energy physics, an effective model is constructed [39]. As shown in Fig. 1(e), the low-lying states at $\Gamma$ are the $|P_{1z}^\uparrow\rangle$ of two Bi layers and $|P_{2z}^\uparrow\rangle$ of two Te layers (Te1 and Te1'), the superscripts $+\uparrow$, $−\downarrow$ stand for parity. The SOC further leads to band inversion. The symmetries of AFM system are the three-fold rotation symmetry $C_{3z}$ and $\mathcal{PT}$. In the basis of $|P_{1z}^\uparrow\rangle, |P_{1z}^\downarrow\rangle, |P_{2z}^\uparrow\rangle, |P_{2z}^\downarrow\rangle$, the representation of the symmetry operations is given by $C_{3z} = \exp[i(\pi/3)\sigma^z \otimes 1]$ and $\mathcal{PT} = i\sigma^y \mathcal{K} \otimes \tau^z$ ($\mathcal{K} = i\sigma^y \mathcal{K} \otimes 1$, $\mathcal{P} = 1 \otimes \tau^z$), where $\mathcal{K}$ is complex conjugation operator, $\sigma^{x,y,z}$ and $\tau^{x,y,z}$ denote the Pauli matrices in the spin and orbital space, respectively. The generic form of the AFM Hamiltonian obeying these symmetries is

$$\mathcal{H}_{\text{AFM}}(\mathbf{k}) = A_1 k_z \sigma^z \otimes \tau^z + A_2 (k_x \sigma^x - k_y \sigma^y) \otimes \tau^z + M_4 (\mathbf{k}) 1 \otimes \tau^z + M_5 (\mathbf{k}) 1 \otimes \tau^y + \epsilon_0 (\mathbf{k}), \quad (1)$$

where $\epsilon_0 (\mathbf{k}) = C_0 + C_1 k_x^2 + C_2 (k_x^2 + k_y^2)$, $M_4 (\mathbf{k}) = M_4 + M_4 k_z^2 + M_2 (k_x^2 + k_y^2)$, $M_5 (\mathbf{k}) = B_0 + B_1 k_z^2 + B_2 (k_x^2 + k_y^2)$, and $A_2 (k_z) = A_2 + A_3 k_z$. Here $M_0 < 0$ and $M_1, M_2 > 0$ correctly characterizes the band inversion at $\Gamma$ [42]. Without $M_5 (\mathbf{k})$ and $A_3$ terms, Eq. (1) is nothing but the textbook TI model in Bi$_2$Te$_3$ family with a single surface Dirac cone [43]. $M_5 (\mathbf{k})$ and $A_3$ are $\mathcal{P}$, $\mathcal{T}$-breaking perturbations induced by the $z$-direction Néel order on Mn. The direct consequence of $M_5 (\mathbf{k})$ term is to open a gap in the surface-state spectrum with the sign independent of the surface orientation, which is equivalent to induce a hedgehodge magnetization on the TI surface. This is confirmed by the (110) and (001) surface spectra by first-principles calculations in Fig. 1(f), which are different from the gapless Dirac state on $\mathcal{T} \tau_{1/2}$-preserving surface in MBT. Also, they are different from the $x$ axis AFM state in M$_2$BT, which has gapless surface state on the surfaces parallel to Néel order with surface Dirac cone shifted away from $\Gamma$ [42]. The hedgehodge-like surface gap of AFM-$z$ M$_2$BT make it an ideal platform for topological magnetoelectric effect.

The AFM ground state of M$_2$BT could be tuned into the FM state by a magnetic field. From above we see low energy physics in AFM M$_2$BT is described by a TI model and $\mathcal{T}$, $\mathcal{P}$-breaking perturbations. For a $z$-axis FM order,

| Thickness (NL) | $\Delta E_{\text{AFM}}$ (meV/Mn) | MAE (meV/Mn) | $E_{\gamma}$ (AFM) (meV) | $E_{\gamma}$ (FM) (meV) |
|----------------|---------------------------------|--------------|-------------------------|-------------------------|
| 1              | $-3.726$                        | 0.052        | 407.9                   | 83.6                    |
| 2              | $-3.953$                        | 0.094        | 67.1                    | 43.1                    |
| 3              | $-3.973$                        | 0.108        | 29.5                    | 60.7                    |
| 4              | $-4.031$                        | 0.104        | 31.0                    | 52.8                    |
| 5              | $-4.001$                        | 0.109        | 20.0                    | 41.2                    |
| 6              | $-4.098$                        | 0.102        | 24.5                    | 31.1                    |
| 7              | $-4.028$                        | 0.103        | 19.4                    | 19.3                    |
| $\infty$ (bulk)| $-3.434$                        | 0.117        | 50.6                    | 0                       |

TABLE I. Thickness dependence of M$_2$BT films magnetism and the energy gap in AFM and FM states. $\Delta E_{\text{AFM}} = E_{\text{FM}} - E_{\text{AFM}}$ is the total energy difference of the AFM and FM states along $z$ direction. The Néel type AFM is likely the ground state. The AFM film are AxI with $\mathcal{C} = 0$; while for the FM films, $\mathcal{C} = 0$ for 1 NL and $\mathcal{C} = 1$ for 2-7 NL from first-principles calculations.
a $T$-breaking but $P$-conversing perturbation should be added, and the resulting possible phase is either Weyl semimetal, 3D QAH or trivial magnetic insulator [13, 44]. The band structure of $z$-axis FM $M_2$BT bulk in Fig. 2 displays a pair of band crossings at Weyl points (W and W) along the $\Gamma$-A line. The Wilson loop calculations suggest the Chern number $C = 1$ at $k_z = 0$ and $C = 0$ at $k_z = \pi$ plane [42], which is consist with the minimal ideal Weyl semimetal in Fig. 2(c). Figs. 2(d)-(f) shows surface-state on different typical surfaces, where Fermi arcs on (1\bar{1}0) termination are clearly seen in Fig. 2(e). Explicitly, the additional $T$-breaking but $P$-conversing terms describing the $z$-axis FM state is

$$
\delta H_{FM}(\mathbf{k}) = A_1^z k_z \sigma^x + A_2^z (k_y \sigma^x - k_x \sigma^y) \otimes \tau^y + M_1^z(\mathbf{k}) \sigma^z \otimes 1 + M_2^z(\mathbf{k}) \sigma^z \otimes \tau^z,
$$

where $M_j^z(\mathbf{k}) = D_0^j + D_1^j k_x^2 + D_2^j (k_x^2 + k_y^2)$ with $j = 1, 2$. This model is similar to FM MBT but with different parameters [42].

Now we understand that the magnetic TI $M_2$BT is well described by a $\text{Bi}_2\text{Te}_3$-type TI model with corresponding $T$-breaking perturbations introduced by Mn. The band inversion in 3D suggests the nontrivial topology may also exist in 2D, which we characterize in below. FM $M_2$BT films have $PT$ symmetry which leads to $C = 0$. They are $A\chi I$ with $t$ nonquantized from $P$, $T$ breaking and finite-size effect. We calculate the energy level versus the film thickness of AFM Hamiltonian (1). Due to quantum confinement, the bulk bands become 2D subbands. As shown in Fig. 3(a), the gap converges when the film exceeds 5 NL, which is consistent with the first-principles calculations listed in Table I. The first pair of subbands $|S1\rangle$ and $|S2\rangle$ are localized on the two surfaces of the thin film [42, 45], with a decay length of about 2 NL. For FM films, $T$ and $PT$-breaking but $P$-conversing leads to spin polarized bands, allowing $C \neq 0$. As calculated in Table I, $C = 0$ for 1 NL and $C = 1$ for 2-7 NL. The spin polarized energy level versus the film thickness in the FM state is calculated in Fig. 3(b), and $C$ is determined by the number of polarized band inversion [46]. Fig. 3(b) suggests 3 NL has the maximum gap in $C = 1$ QAH and is consistent with first-principles calculations.

Intriguingly, here as the interlayer exchange coupling is quite weak, the Mn layers may be driven into different magnetic configurations, which further modify the band topology. Take 2 NL for example, we calculate the band structure, relative total energy and $C$ for five different magnetic configurations named AFM, FM, interstate I, II, III shown in Fig. 4. Clearly, FM, I and II are QAH with $C = 1$. AFM and III have $C = 0$, lead to zero Hall conductance. As we show below, the AFM state is $A\chi I$ but III is trivial insulator.

To describe the layer magnetization tuned QAH state in $M_2$BT film, we start with the low energy physics which is well described by the massive Dirac surface states only, where the intrinsic magnetic ordering introduces different Zeeman terms on these two surfaces. The generic effective Hamiltonian for thin film is

$$
H_{\text{film}}(k_x, k_y) = v_F (k_y \sigma_x - k_x \sigma_y) \otimes \tau_z + m(k) 1 \otimes \tau_x + g_a \sigma_z \otimes \tau_z + g_f \sigma_z \otimes 1,
$$

with the basis of $|\uparrow\rangle$, $|\downarrow\rangle$, $|b\uparrow\rangle$, and $|b\downarrow\rangle$, where $t$ and $b$ denote the top and bottom surface states, and $\uparrow$ and $\downarrow$ represent spin up and down states, respectively. $\sigma_i$ and $\tau_i (i = x, y, z)$ are Pauli matrices acting on spin and layer, respectively. $v_F$ is the Fermi velocity. $m(k) = m_0 + m_1 (k_x^2 + k_y^2)$ is the hybridization between the top and bottom surface states. The third and fourth terms describe the Zeeman-type spin splitting of top $g_t$ and bottom $g_b$ surface states induced by the FM exchange
of Mn along z axis, where $g_{a,f} = (g_t \mp g_b)/2$ are the staggered and uniform Zeeman field, respectively. Both $\uparrow$ and $\downarrow$ FM Mn layers will contribute to the Zeeman field. In the mean field approximation,

$$g_i = \sum_j \text{sgn}(s_j^z) \lambda^{ij}, \quad (i = t, b)$$

where $j$ labels the Mn layer index, $s_j^z$ is the z-component of Mn local spin in layer $j$, $\lambda^{ij}$ is the effective exchange parameter between local moments in layer $j$ and the top ($i = t$) or bottom ($i = b$) surface states, respectively. All $\lambda^{ij}$ have the same sign and we set $\lambda_{ij} > 0$. $\text{sgn}(s_j^z)$ comes from the magnetization direction of each Mn layer.

The Hamiltonian (3) describes both QAH and AxI states characterized by $C$. The band dispersion is given by

$$\epsilon_j^2(k_x, k_y) = v_F^2(k_x^2 + k_y^2) + (\sqrt{m(k_x^2 + g_{ij}^2)} \pm g_f)$$

$C$ only changes at the gap closing point determined by $\sqrt{m_F^2 + g_{ij}^2} = |g_f|$. When $\sqrt{m_F^2 + g_{ij}^2} < |g_f|$, the system is QAH with $C = g_f/|g_f|$; when $\sqrt{m_F^2 + g_{ij}^2} > |g_f|$, the system has $C = 0$. AFM and III are topologically equivalent but have quite different origins. In AFM, with opposite magnetic exchange coupling on two surfaces, $g_f = 0$ and $g_a$ is finite, it is AxI with a nonquantized but finite $\theta$ response [47]. While in III, $g_a = 0$ and the hybridization gap exceeds the finite FM exchange gap $g_a$, thus it is a trivial insulator with $\theta = 0$.

With the gap and $C$ in these magnetic states, $\lambda^{ij}$ can be determined. For 2 NL, we approximate $\lambda_1 \equiv \lambda_{11} = \lambda^{b1}$, $\lambda_2 \equiv \lambda_{12} = \lambda^{b2}$, $\lambda_3 \equiv \lambda_{13} = \lambda^{b3}$, and $\lambda_4 \equiv \lambda_{14} = \lambda^{b4}$. This yields $\lambda_1 = 19.0$ meV, $\lambda_2 = 2.3$ meV, $\lambda_3 = 18.4$ meV, $\lambda_4 = 1.8$ meV and $m_0 = 3.3$ meV. Consistent with 2 NL decay length of surface states, $\lambda^{ij}$ almost vanishes when $i = t$ and $j \geq 5$ as determined from the magnetic states in 3 NL [42].

Finally, we discuss the field-induced magnetic transitions. The evolution of magnetic transition under external field can be described by a magnetic bilayer Stoner-Wohlfarth model with an interlayer exchange coupling $J_{1,2}$ and an effective anisotropy $K$. From Fig. 4, within each NL the interlayer AFM coupling is $J_1 = 0.8$ meV, and $J_2 = 0.35$ meV between adjacent NL. The uniaxial anisotropy $K \approx 0.1$ meV in Table I. With the field applied parallel to the magnetic easy z axis, since $K \ll J_{1,2}$, the AFM ground state undergoes a spin-flop transition to a canted state where the sublattice magnetization is roughly perpendicular to z axis. Further increasing the field brings the canted magnetizations to FM state by coherent rotation. Take 2 NL as an example, the z-axis magnetic field will drive AFM to a canted state at field $H^{ci}_z$, then to II, and finally to FM. For II, it is energetically favorable than I. A rough estimation yields $H^{ci}_z \approx 1.6$ T. The coherent rotation of layer magnetization and the corresponding QAH plateau transition (with $C = 0 \rightarrow 1$) at small fields provides a promising platform for chiral Majorana fermion.

The intrinsic van der Waals magnetic material M2BT hosts rich topological quantum states in different spatial dimensions, which is well described by a Bi$_2$Te$_3$-type TI model with certain $T$-breaking perturbations. We expect superlattice-like new magnetic TI such as M2BT/MBT and M2BT/Bi$_2$Te$_3$ with tunable exchange interactions.
and topological properties may be fabricated. This will further enrich the magnetic TI family and provide a new material platform for exotic topological phenomena.

We acknowledge Z. Liu, H. Zhang, W. Li, G. Xu, H. Yang, J. Zhu and Y. Zhang for valuable discussions. This work is supported by the Natural Science Foundation of China through Grant Nos. 11774055, 11574051, and 11874117, the National Key Research Program of China under Grant Nos. 2016YFA0300703, 2019YFA0308404, the Natural Science Foundation of Shanghai under Grant No. 19ZR1471400. Y.L. and Y.J. contributed equally to this work.

* wjingphys@fudan.edu.cn

[1] M. Z. Hasan and C. L. Kane, Rev. Mod. Phys. 82, 3045 (2010).
[2] X.-L. Qi and S.-C. Zhang, Rev. Mod. Phys. 83, 1057 (2011).
[3] Y. Tokura, K. Yasuda, and A. Tsukazaki, Nat. Rev. Phys. 1, 126 (2019).
[4] J. Wang and S.-C. Zhang, Nature Mat. 16, 1062 (2017).
[5] F. D. M. Haldane, Phys. Rev. Lett. 61, 2015 (1988).
[6] X.-L. Qi, Y.-S. Wu, and S.-C. Zhang, Phys. Rev. B 74, 085308 (2006).
[7] R. Yu, W. Zhang, H.-J. Zhang, S.-C. Zhang, X. Dai, and Z. Fang, Science 329, 61 (2010).
[8] X.-L. Qi, T. L. Hughes, and S.-C. Zhang, Phys. Rev. B 78, 195424 (2008).
[9] A. M. Essin, J. E. Moore, and D. Vanderbilt, Phys. Rev. Lett. 102, 146805 (2009).
[10] R. Li, J. Wang, X. L. Qi, and S. C. Zhang, Nature Phys. 6, 284 (2010).
[11] J. Wang, B. Lian, X.-L. Qi, and S.-C. Zhang, Phys. Rev. B 92, 081107 (2015).
[12] T. Morimoto, A. Furusaki, and N. Nagaosa, Phys. Rev. B 92, 085113 (2015).
[13] J. Wang, B. Lian, and S.-C. Zhang, Phys. Rev. B 93, 045115 (2016).
[14] M. Mogi, M. Kawamura, R. Yoshimi, A. Tsukazaki, Y. Kozuka, N. Shirakawa, K. S. Takahashi, M. Kawasaki, and Y. Tokura, Nature Mater. 16, 516 (2017).
[15] N. Read and D. Green, Phys. Rev. B 61, 10267 (2000).
[16] J. A. Alica, Rep. Prog. Phys. 75, 076501 (2012).
[17] X.-L. Qi, T. L. Hughes, and S.-C. Zhang, Phys. Rev. B 82, 184516 (2010).
[18] A. Kitaev, Ann. Phys. 303, 2 (2003).
[19] C. Nayak, S. H. Simon, A. Stern, M. Freedman, and S. Das Sarma, Rev. Mod. Phys. 80, 1083 (2008).
[20] B. Lian, X.-Q. Sun, A. Vaezi, X.-L. Qi, and S.-C. Zhang, Proc. Natl. Acad. Sci. U.S.A. 115, 10938 (2018).
[21] C.-Z. Chang, J. Zhang, X. Feng, J. Shen, Z. Zhang, M. Guo, K. Li, Y. Ou, P. Wei, L.-L. Wang, Z.-Q. Ji, Y. Feng, S. Ji, X. Chen, J. Jia, X. Dai, Z. Fang, S.-C. Zhang, K. He, Y. Wang, L. Lu, X.-C. Ma, and Q.-K. Xue, Science 340, 167 (2013).
[22] J. G. Checkelsky, R. Yoshimi, A. Tsukazaki, K. S. Takahashi, Y. Kozuka, J. Falson, M. Kawasaki, and Y. Tokura, Nature Phys. 10, 731 (2014).
[23] A. J. Bestwick, E. J. Fox, X. Kou, L. Pan, K. L. Wang, and D. Goldhaber-Gordon, Phys. Rev. Lett. 114, 187201 (2015).
[24] C.-Z. Chang, W. Zhao, D. Y. Kim, H. Zhang, B. A. Assaf, D. Heiman, S.-C. Zhang, C. Liu, M. H. W. Chan, and J. S. Moodera, Nature Mater. 14, 473 (2015).
[25] L. He, K. L. Billinge, R. Zhong, J. A. Schneeloch, T. Liu, T. Valla, J. M. Tranquada, G. Gu, and J. C. S. Davis, Proc. Natl. Acad. Sci. USA 112, 1316 (2015).
[26] Q. L. He, L. Pan, A. L. Stern, E. C. Burks, X. Che, G. Yin, J. Wang, B. Lian, Q. Zhou, E. S. Choi, K. Murata, X. Kou, Z. Chen, T. Nie, Q. Shao, Y. Fan, S.-C. Zhang, K. Liu, J. Xia, and K. L. Wang, Science 357, 294 (2017).
[27] M. Kayyalha, D. Xiao, R. Zhang, J. Shin, J. Jiang, F. Wang, Y.-F. Zhao, R. Xiao, L. Zhang, K. M. Fjelkkowski, P. Maudal, M. Winnerlein, C. Gould, Q. Li, L. W. Molenkamp, M. H. W. Chan, N. Samarth, and C.-Z. Chang, Science 367, 64 (2020).
[28] J. Wang, Q. Zhou, B. Lian, and S.-C. Zhang, Phys. Rev. B 92, 064520 (2015).
[29] W. Ji and X.-G. Wen, Phys. Rev. Lett. 120, 107002 (2018).
[30] Y. Huang, F. Setiawan, and J. D. Sau, Phys. Rev. B 97, 100501 (2018).
[31] B. Lian, J. Wang, X.-Q. Sun, A. Vaezi, and S.-C. Zhang, Phys. Rev. B 97, 125408 (2018).
[32] D. Zhang, M. Shi, T. Zhu, D. Xing, H. Zhang, and J. Wang, Phys. Rev. Lett. 122, 206401 (2019).
[33] J. Li, Y. Li, S. Du, Z. Wang, B.-L. Gu, S.-C. Zhang, K. He, W. Duan, and Y. Xu, Sci. Adv. 5, eaaw5685 (2019).
[34] Y. Gong, J. Guo, J. Li, K. Zhu, M. Liao, X. Liu, Q. Zhang, L. Gu, L. Tang, X. Feng, D. Zhang, W. Li, C. Song, L. Wang, P. Yu, X. Chen, Y. Wang, H. Yao, W. Duan, Y. Xu, S.-C. Zhang, X. Ma, Q.-K. Xue, and K. He, Chin. Phys. Lett. 36, 076801 (2019).
[35] M. M. Otrokov, A. Yakimovskikh, H. Bentmann, A. Zeugner, Z. S. Aliev, G. S. A. W. Boller, A. R. V. Koroleva, D. Estyunin, A. M. Shikin, M. Blanco-Rey, M. Hoffmann, A. R. V. Vyazovskaya, S. V. Ereemenko, Y. M. Koroteev, I. R. Amiralsanov, M. B. Babanyan, N. T. Mamedov, N. A. Abdullayev, V. N. Zverev, B. Buckner, E. F. Schwier, S. Kumar, A. Kimura, L. Petaccia, G. Di Santo, R. C. Vidal, S. Schatz, K. Kijhn, C.-H. Min, S. K. Moser, T. R. F. Peixoto, F. Reinert, A. Ernst, P. M. Echenique, A. Isaeva, and E. V. Chulkov, Nature 576, 416 (2019).
[36] Y. Deng, Y. Yu, M. Z. Shi, J. Wang, X. H. Chen, and Y. Zhang, arXiv:1904.11468.
[37] C. Liu, Y. Wang, H. Li, Y. Wu, Y. Li, J. Li, K. He, Y. Xu, J. Zhang, and Y. Wang, Nat. Mater. (2020) doi:10.1038/s41563-019-0573-3.
[38] J. Ge, Y. Liu, J. Li, H. Li, T. Luo, Y. Wu, Y. Xu, and J. Wang, arXiv:1907.09947.
[39] J. Zhang, D. Wang, M. Shi, T. Zhu, H. Zhang, and J. Wang, arXiv e-prints , arXiv:1906.07891 (2019), arXiv:1906.07891 [cond-mat.mes-hall].
[40] Y. Lv et al., to be published (2020).
[41] Y. Hou and R. Wu, Nano Lett. 19, 2472 (2019).
[42] See Supplemental Material for technical details.
[43] H. Zhang, C.-X. Liu, X.-L. Qi, X. Dai, Z. Fang, and S.-C. Zhang, Nature Phys. 5, 438 (2009).
[44] A. A. Burkov and L. Balents, Phys. Rev. Lett. 107,
[45] C.-X. Liu, H.-J. Zhang, B. Yan, X.-L. Qi, T. Frauenheim, X. Dai, Z. Fang, and S.-C. Zhang, Phys. Rev. B 81, 041307 (2010).

[46] J. Wang, B. Lian, H. Zhang, Y. Xu, and S.-C. Zhang, Phys. Rev. Lett. 111, 136801 (2013).

[47] Z. Liu and J. Wang, arXiv:1912.07786 (2019).