Gapped Surface States in a Strong-Topological-Semimetal

A.P. Weber, Q.D. Gibson, Huiwen Ji, A.N. Caruso, A.V. Fedorov, R.J. Cava, and T. Valla

Department of Physics and Astronomy, University of Missouri-Kansas City, Kansas City, Missouri 64110, USA
Department of Chemistry, Princeton University, Princeton, New Jersey 08544, USA
Advanced Light Source, Lawrence Berkeley National Laboratory, Berkeley, California 94720, USA
Condensed Matter Physics and Materials Science Department, Brookhaven National Laboratory, Upton, New York 11973, USA

A three-dimensional strong-topological-insulator or -semimetal hosts topological surface states which are often said to be gapless so long as time-reversal symmetry is preserved. This narrative can be mistaken when surface state degeneracies occur away from time-reversal-invariant momenta. The mirror-invariance of the system then becomes essential in protecting the existence of a surface Fermi surface. Here we show that such a case exists in the strong-topological-semimetal Bi$_2$Se$_3$. Angle-resolved photoemission spectroscopy and ab initio calculations reveal partial gapping of surface bands on the Bi$_2$Se$_3$-termination of Bi$_2$Se$_3$(111), where an 85 meV gap along $\Gamma K$ closes to zero toward the mirror-invariant $\Gamma M$ azimuth. The gap opening is attributed to an interband spin-orbit interaction that mixes states of opposite spin-helicity.

PACS numbers: 71.20.-b, 73.20.At, 79.60.Bm, 71.70.Ej

Topological insulator materials (TIM) [1], which include insulators [2–4] and semimetals [5–10], possess an inverted bulk band gap that hosts unusually robust, spin-helical topological surface states (TSS) at the material’s boundaries. The presence of TSS is guaranteed by the topology of the bulk bands and the states can only be removed from the Fermi energy if protecting symmetries are broken. Strong topological insulator (STI) [11, 12] materials hold a special distinction, as they are often said to host gapless TSS protected by time-reversal symmetry (TRS) on every surface termination of the crystal. We will show it is possible for the TSS of material in a STI phase to intrinsically acquire a finite gap at all momenta in the surface Brillouin zone (SBZ) which do not lie on a mirror-invariant azimuth. In that case, the mirror-symmetry (MS) of the crystal lattice must be intact to guarantee the existence of a surface Fermi surface. Earlier work by Teo, Fu, and Kane [13] had foreseen this possibility, however, no concrete examples have been obtained in experiment. Here we show that such a system exists in the strong-topological-semimetal Bi$_2$Se$_3$. The finding of mirror-protected TSS in this system also provides direct confirmation that the strong-topological-insulator phase can coexist with the topological-crystalline-insulator phase of matter, as was suggested by Rauch et al. [14] for the case of Bi-chalcogenides.

Bi$_2$Se$_3$ is a rhombohedral superlattice material consisting of alternating Bi$_2$ and Bi$_2$Se$_3$ layers stacked along the (111) direction. Previously, Dirac cone-like TSS were found in the gap between the first fully occupied bulk valence band (BVB) and the hole-like bulk conduction band (HBCB) on two different surface terminations of Bi$_2$Se$_3$(111) [10]. It was determined that the TSS result from a parity inversion at the $\Gamma$-point of the bulk Brillouin zone (BBZ), a characteristic shared with the Bi$_2$Se$_3$ parent compound [15]. Although effects of hybridization between surface and bulk electrons were mentioned, it was not made clear that the states of the upper Dirac cone appearing on the Bi$_2$Se$_3$-terminated surface, which have an electron-like dispersion, must have crossed the HBCB to reach the Fermi level. The possibility of new topological constraints on the surface electron-structure above the HBCB was also not explored. Earlier work [9] predicted that the band gap above the HBCB is characterized by a single parity inversion at the F-point of the BBZ, which demands that TSS within this gap must come in an odd number of pairs [13]. If that is the case, then an odd number of surface bands must appear within the gap to meet the state which crossed the HBCB. Moreover, this parity inversion is away from the center of the BBZ and the crystal does not cleave at the center of inversion. Under these conditions, the TSS pairs are not constrained to have Dirac points at time-reversal invariant momenta (TRIM) [15].

Consistent with the prediction in ref. [9], it is in the momentum-space region outside the HBCB edge that evidence of partially gapped TSS is found through ab initio calculations and angle-resolved photoemission spectroscopy (ARPES) measurements on the Bi$_2$Se$_3$-terminated surface of Bi$_2$Se$_3$(111). The TSS degenerate away from TRIM on the mirror-invariant $\Gamma M$ azimuth in the SBZ. The bands become separated elsewhere, and an 85 meV gap between the surface state branches is measured along $\Gamma K$ in ARPES. The origin of this gap is accounted for in a model for the spin-orbit interaction on a (111) crystal surface. These findings place 2D electron hybridization within the subject of pristine STI materials and provide direct evidence for MS protection of surface states in a Bi-chalcogenide.

Single crystals of Bi$_2$Se$_3$ were synthesized following a previously reported procedure [10]. ARPES was per-
formed using a Scienta SES-100 electron spectrometer at beamline 12.0.1 of the Advanced Light Source with a combined instrumental energy resolution of ~12 meV and an angular resolution better than ± 0.07 Å. The sample was cleaved under ultrahigh vacuum conditions (< 5.0 x 10^{-9} Pa) and kept at ~15 K. Temperature was measured using a silicon sensor mounted near the sample.

Electron-structure calculations were performed in the framework of density functional theory (DFT) using the WIEN2K code [16] with a full-potential linearized augmented plane-wave and local orbitals basis together with the Perdew-Burke-Ernzerhof [17] parametrization of the generalized gradient approximation, using a slab geometry. Experimentally determined lattice parameters and atom positions were used to construct the slabs. The plane-wave cutoff parameter RMTKmax was set to 7 and the Brillouin zone was sampled by 9 k points. Spin-orbit coupling (SOC) was included. To study the Bi2Se3-terminated surface, a slab was constructed of 6 Bi2Se3 layers and 5 Bi2 layers, with 10 Å of vacuum between adjacent slabs. The contribution of the surface atoms to the overall surface electronic structure was determined by calculating the partial contribution of each atomic basis set to the wave functions at all k points.

The calculated electron-structure for Bi2Se3-terminated slabs of Bi2Se3(111) is shown in Figure 1. A surface state crossing protected by MS only is indicated by the red circle and the blue circle indicates a crossing which is "dually protected" [14] by both MS and TRS. The TSS within the blue-circled region lie between the BVB and HBCB. These TSS were the primary focus of previous investigations [9, 10]. The TSS that cross within the red-circled region, which are the focus of the present work, lie above the HBCB. The parity invariants of the bulk band structure counted up to the HBCB were previously determined to be +1, +1, +1, and −1 at the Γ, Z, L, and Φ points of the BZ, respectively [9]. The product of the parity invariants is −1, which characterizes the gap above the HBCB as a STI-type [11, 12]. Applying the methods of ref. [13] to our case, an odd number of TSS pairs are expected to exist between the surface projections of F (M) and Γ (Γ). Indeed, we observe a single pair of TSS that cross each other along the ΓM azimuth and degenerate with different ends of the bulk gap at Γ and M. Group-theoretical considerations indicate why this crossing is allowed even while the surface states are seen to be gapped along the ΓΦ azimuth.

At the (111) surface, the R3m symmetry of the crystal reduces to C3v. For wave-vectors lying between Γ and M, the point-group symmetry reduces to C3, which contains two irreducible representations characterized by mirror eigenvalues of ±i. Through the definition of the mirror operation [13], it is easily shown that the two irreducible representations correspond to states of opposite spin-helicity, which cannot hybridize with each other on the mirror-invariant ΓM azimuth. This explains why the crossing circled in red is allowed and, indeed, protected by the crystal’s mirror symmetry. In contrast, the point group of the wave-vectors along ΓΦ is C1. By symmetry, crossings between Γ and Φ are avoided, even for states of opposite spin-helicity. The same is true for all wave-vectors in the SBZ which do not lie on a mirror-invariant azimuth. The double-arrow in panel (a) indicates what can therefore be understood as a hybridization gap Δ resulting from the avoided crossing of spin-helical surface states, as will be discussed later. Note that the combination of C3 and time-reversal symmetry imply that the surface Fermi surface will consist of six equivalent pockets that each enclose a surface state degeneracy point. This observation is consistent with the definition of a STI material as put forward by Teo, Fu, and Kane [13].

The k-dependence of the gapped structure results from competing SOC interactions which couple to different components of the spin degree of freedom. This can be captured in a model for the SOC Hamiltonian $H_{soc} \propto (\mathbf{\sigma} \times \nabla V) \cdot \mathbf{\sigma}$ using $k \cdot p$ theory. Polar coordinates are chosen with \(\hat{\Gamma}\) as the origin and θ denoting...
the in-plane azimuthal angle from the \( k_z \)-axis, aligned to the \( \Gamma K \) direction. If the energy-splitting of the TSS near \( k = 0 \) is taken to equal \( 2\kappa v \), then \( H_{soc} \) to first-order in \( k \) couples the in-plane, tangential component of spin \( \langle \sigma_r \rangle \) to the out-of-plane electrostatic potential gradient with a strength \( v \) as \( H_1 \left( \frac{k}{\kappa} \right) \equiv v(k - \kappa)\sigma_r \). To third-order in \( k \), there appears a second term \([18]\) that couples the out-of-plane component of spin to the in-plane crystalline potential gradient with a strength \( \lambda \) as \( H_2 \left( \frac{k}{\kappa} \right) \equiv \lambda k^3\cos(3\theta)\sigma_z \). Together,

\[
H_{soc} \left( \frac{k}{\kappa} \right) = \begin{bmatrix}
\lambda k^3\cos(3\theta) & iv(k - \kappa)e^{-i\theta} \\
-iv(k - \kappa)e^{i\theta} & -\lambda k^3\cos(3\theta)
\end{bmatrix}
\]

and we find the spin-orbit contribution to the dispersion

\[
E_{soc}(k) = \pm \sqrt{v^2(k - \kappa)^2 + \lambda^2 k^6\cos^2(3\theta)}
\]

where we refer to + and − as the upper Dirac branch (UDB) and lower Dirac Branch (LDB), respectively. At \( k = \kappa \), the magnitude of the gap between the DBs is determined solely by the second term under the square root. The gap carries the sign of an \( f \)-wave, which changes at the mirror-invariant azimuths, signifying a change in the \( z \)-polarizations of the DBs (this could alternately be described as a crossing of bands with positive and negative \( z \)-polarization). The sign appears in the interband matrix element

\[
\Delta = \left\langle -H_2 \left( \frac{k}{\kappa} \right) \right| + = -2\lambda k^3\cos(3\theta)
\]

for the spin-helical states typically invoked in the discussion of simple, gapless TSS. In this sense, the crystalline anisotropy should cause an interband SOC effect that gives rise to the gapped structure. The model only differs from the TSS of the STI Bi\(_2\)Te\(_3\) \([18]\) in the location of the Dirac point. No extra bands need to be inferred to achieve the complex, partially gapped TSS described below.

Fig. 2 displays electron-structure calculated from the model Hamiltonian, for which we have chosen the values \( v = 5.5 \text{ eVÅ}, \lambda = 55 \text{ eVÅ}^3, \kappa = 0.25 \text{ Å}^{-1} \) in rough approximation of the Bi\(_2\)Se\(_3\)-termination electron-structure. The reader should note the many omissions from this model such as spin-orbital entanglement (which will disallow a pure spin-eigenstate character for the TSS and limit the magnitude of spin-polarization) \([19]\), and higher-order interaction terms \([20]\). Fig. 2(a) shows the spin-resolved band structure along the \( \theta = 0 \) azimuth, with the bands corresponding to the case \( \lambda = 0 \) plotted in black. The minimum energy gap \( \delta \) is located inside of \( k = \kappa \), and we find it is the case that at \( k_3 \leq \kappa \) for all \( \theta \), as shown in panel (c). It is telling to inspect the in-plane helical spin-polarization, indicated by color scale, relative to the magnitude of the out-of-plane polarization, indicated by marker size. For \( k < 0.1 \text{ Å}^{-1} \), the spin is helical, but as the contribution of \( H_2 \) grows, the deviation from linear dispersion is accompanied by increasing \( z \)-polarization. At \( k_3 \) the contribution of \( H_2 \) overtakes that of \( H_1 \). At \( k = \kappa \), the \( z \)-polarization is 100% and then attenuates to \( \sim 90\% \) for \( k > \kappa \), where the spin-helicity has reversed. The scenario of competing SOC interactions takes place throughout the SBZ, resulting in unusual constant energy contour (CEC) shapes and topologies, displayed in panel (b). If the electron-dispersion were determined solely by the SOC in this model, there would exist two Lifshitz points \([21]\) in the chemical potential located at \( \mu = \pm E(k_3, \theta = 0) \). At these points, the Fermi surface topology changes from six pockets enclosing the TSS degeneracies to one electron-pocket and one hole-pocket enclosing \( \Gamma \). Near a Lifshitz point, straight edges in the CECs are centered on the \( \Gamma M \) direction, opposite of what would be expected for simple, gapless TSS on a (111) surface \([18]\). This same pattern appears in CECs probed by ARPES, described in Fig.3(c) below.

![Figure 2](image-url)
Figure 3. ARPES electron-structure near the center of the 1st surface Brillouin zone: The Fermi surface (a,b), and constant energy contour at -185 meV (c). Band structure parallel to ΓM direction for \( k_x = 0 \) (d) and \( k_x = 0.21 \) Å\(^{-1}\) (i). Band structure parallel to ΓK direction at \( k_y = -0.255 \) Å\(^{-1}\) (e), \( k_y = 0.255 \) Å\(^{-1}\) (f), and \( k_y = 0 \) (h). Energy distribution curves at the degeneracy point \( (k_x, k_y) = (0, -0.255) \) Å\(^{-1}\) (g) and the saddle point of the lower Dirac branch \( (k_x, k_y) = (0.21, 0) \) Å\(^{-1}\) (j). The momentum-space locations for panels (d-f) and (h-i) are indicated by lines overlayed onto the Fermi surface in (a). The gray-scale varies from white-to-black following the minimum-to-maximum photoemission intensity within each image individually.

The observation of gapped surface states on a (111) surface with large SOC is not without precedent, having been previously identified in heterostructures with BiAg\(_2\) surface alloys [23], however, the mechanism for the "interband SOC" between the antiparallel, spin-helical surface states in that case was left unspecified. The two-term model Hamiltonian approach shown in this work could be extended to describe not only BiAg\(_2\) surface states, which are known to have a sizable coupling to the in-plane crystal potential gradient [24], but also many other systems, whether topologically trivial or not, in which spin-helical states intersect away from Kramer’s momenta, such as Pb quantum wells [25]. TIM, rather than topologically trivial materials, may offer more robust platforms for studying this type of spin-gap physics in 2D electron systems. What is lacking at this time is a straightforward and reliable way of predicting if a given STI surface will possess gapped TSS. The present results indicate a need to merge different conceptualizations of the topological insulator (as defined in the abstract sense band topology [26]) in order to determine the conditions necessary for this phenomenon.

Some have pointed out [13, 14, 27] that several well-known TIM possess bulk electron structure that can be characterized as topologically non-trivial using separate methods. The parity invariant method is used to characterize \( Z_2 \) topological insulators [11, 12], which include STIs, while mirror topological crystalline insulators (TCIs) are characterized by a non-zero difference in the number of counterpropagating "edge states" corresponding to a crystallographic mirror plane [13, 28, 29]. Teo, Fu, and Kane [13] had considered that the crossing of TSS at non-TRIM was a possibility for a strong \( Z_2 \) topological insulator surface, which motivated them to develop the foundational theory for TCIs. This Letter has presented an experimentally realized case in which the concepts of \( Z_2 \) topological insulators and TCIs have become entangled beyond precedent; completely breaking the MS would allow the TSS to become fully gapped, even while TRS remains unbroken. Surely, the significance of the \( Z_2 \) topology in determining the electronic physics at all of the possible surfaces of a STI should be revisited.
The financial support of the National Science foundation, Grants No. NSF-DMR-0819860 and No. NSF-DMR-1104612, DARPA-SPAWAR Grant No. N6601-11-1-4110, and the ARO MURI program, Grant No. W911NF-12-1-0461, is gratefully acknowledged. The Advanced Light Source is supported by the US DOE, Office of Basic Energy Sciences, under Contract No.DE-AC02-05CH11231. Brookhaven National Laboratory is supported by the US Department of Energy, Office of Basic Energy Sciences, Contract No. DE-AC02-98CH10886.

*apwnq5@mail.umkc.edu

[1] Y. Ando, J. Phys. Soc. Jpn. **82**, 102001 (2013)
[2] M.Z. Hasan and C.L. Kane, Rev. Mod. Phys. **82**, 3045 (2010)
[3] M. Z. Hasan and J. E. Moore, Annu. Rev. Condens. Matter Phys. **2**, 55 (2011)
[4] X.L. Qi and S.C. Zhang, Rev. Mod. Phys. **83**, 1057 (2011)
[5] D. Hsieh, Y. Xia, L. Wray, D. Qian, A. Pal, J. H. Dil, J. Osterwalder, F. Meier, G. Bihlmayer, C. L. Kane, Y. S. Hor, R. J. Cava, and M. Z. Hasan, Science **323** (2009) 919
[6] D. Hsieh, L Wray, D. Qian, Y. Xia, J.H. Dil, F. Meier, L. Patthey, J. Osterwalder, G. Bihlmayer, Y.S. Hor, R.J. Cava, and M.Z. Hasan, New J. Phys. **12** (2010) 125001
[7] A. Stróżecka, A. Eguren, M. Bianchi, D. Guan, C.H. Voetmann, S. Bao, P. Hofmann, and J.I. Pascual, New J. Phys. **14** (2012) 103026
[8] M. Bianchi, D.Guan, A. Stróżecka, C.H. Voetmann, S.Bao, J.I. Pascual, A. Eguren, and P. Hofmann, Phys. Rev. B **85** (2012) 155431
[9] T. Valla, H. Ji, L.M. Schoop, A.P. Weber, Z.H. Pan, J.T. Sadowski, E. Vescovo, A.V. Fedorov, A.N. Caruso, Q.D. Gibson, L. Múchler, C. Felser, and R.J. Cava, Phys. Rev. B **86**, 241101(R) (2012)
[10] Q.D. Gibson, L.M Schoop, A.P. Weber, H. Ji, S. Nadj-Perge, I.K. Drozdov, H. Beidenkopf, J.T. Sadowski, A. Fedorov, A. Yazdani, T. Valla, and R.J. Cava, Phys. Rev. B **88**, 081108(R) (2013)
[11] L. Fu, C.L. Kane, and E.J. Mele, Phys. Rev. Lett. **98**, 106803 (2007)
[12] L. Fu and C. L. Kane, Phys. Rev. B **76**, 045302 (2007)
[13] J. C. Y. Teo, L. Fu, and C. L. Kane, Phys. Rev. B **78**, 045426 (2008)
[14] T. Rauch, M. Flieger, J. Henk, I. Mertig, and A. Ernst, Phys. Rev. Lett. **112**, 016802 (2014)
[15] H. Zhang, C.-X. Liu, X.-L. Qi, X. Dai, Z. Fang, and S.-C. Zhang, Nature Phys. **5**, 438 (2009)
[16] P. Blaha, K. Schwarz, P. Sorantin, and S. Trickey, Comput. Phys. Commun. **59**, 399 (1990)
[17] J. P. Perdew, K. Burke, and M. Ernzerhof, Phys. Rev. Lett. 77, 3865 (1996).
[18] L. Fu, Phys. Rev. Lett. **103**, 266801 (2009)
[19] O.V. Yazyev, J.E. Moore, and S.G. Louie, Phys. Rev. Lett. **105**, 266806 (2010)
[20] S. Basak, H. Lin, L. A. Wray, S.-Y. Xu, L. Fu, M. Z. Hasan, and A. Bansil, Phys. Rev. B **84**, 121401(R) (2011)
[21] I.M. Lifshitz, Sov. Phys. JETP **11**, 1130 (1960)
[22] L. Van Hove, Phys. Rev. **89**, 1189 (1953)
[23] H. Bentmann, S. Abdelouahed, M. Mulazzi, J. Henk, F. Reimert, Phys. Rev. Lett. **108**, 196801 (2012)
[24] C.R. Ast, J. Henk, A. Ernst, L. Moreschini, M.C. Falub, D. Faciél, P. Bruno, K. Kern, and M. Grioni, Phys. Rev. Lett. **98**, 186807
[25] B. Slomski, G. Landolt, S. Muff, F. Meier, J. Osterwalder, and J.H. Dil, New J. Phys. **15**, 125031 (2013)
[26] Any material with a band gap that remains open at every point in the Brillouin zone can be regarded as an insulator in the abstractions of topological band theory. The band structure of a semimetal meeting this criterion can be "smoothly deformed" into that of an insulator.
[27] Y.J. Wang, W.-F. Tsai, H. Lin, S.-Y. Xu, M. Neupane, M. Z. Hasan, and A. Bansil, Phys. Rev. B **87**, 235317 (2013)
[28] L. Fu, Phys. Rev. Lett. **106**, 106802 (2011)
[29] T.H. Hsieh, H. Lin, J. Liu, W. Duan, A. Bansil, and Liang Fu, Nature Comm. **3**, 982 (2012)