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Bahadur Singh, Xiaoting Zhou, Hsin Lin, and Arun Bansil

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# Saddle-like topological surface states on the TT'X family of compounds (T, T' = Transition metal, X= Si, Ge)

Bahadur Singh<sup>b,1,2,\*</sup> Xiaoting Zhou<sup>b,1,2</sup> Hsin Lin,<sup>1,2,†</sup> and Arun Bansil<sup>3</sup>

<sup>1</sup>*Centre for Advanced 2D Materials and Graphene Research Centre,  
National University of Singapore, Singapore 117546*

<sup>2</sup>*Department of Physics, National University of Singapore, Singapore 117542*

<sup>3</sup>*Department of Physics, Northeastern University,  
Boston, Massachusetts 02115, USA*

## Abstract

Topological nodal-line semimetals are exotic conductors that host symmetry-protected conducting nodal lines in their bulk electronic spectrum and nontrivial drumhead states on the surface. Based on first-principles calculations and an effective model analysis, we identify the presence of topological nodal-line semimetal states in the low crystalline symmetric TT'X family of compounds (T, T' = transition metal, X= Si, or Ge) in the absence of spin-orbit coupling (SOC). Taking ZrPtGe as an exemplar system, we show that owing to small lattice symmetry this material harbors a single nodal line on the  $k_y = 0$  plane with large energy dispersion and unique drumhead surface state with a saddle-like energy dispersion. When the SOC is included, the nodal line gaps out and the system transitions to a strong topological insulator state with  $Z_2 = (1; 000)$ . The topological surface state evolves from the drumhead surface state via the sharing of its saddle-like energy dispersion within the bulk energy gap. These features differ remarkably from those of the currently known topological surface states in topological insulators such as Bi<sub>2</sub>Se<sub>3</sub> with Dirac-cone-like energy dispersions.

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<sup>b</sup> These authors contributed equally to this work.

## I. INTRODUCTION

Recent discovery of non-trivial band structures in semimetals has prompted prodigious research interest in topological semimetals<sup>1–17</sup>. Unlike topological insulators (TIs), whose protected gapless states live only on their surfaces, topological semimetals (TSMs) feature unusual crystal-symmetry-protected states both in the bulk as well on the surface. In particular, the bulk Fermi surfaces of TSMs enclose non-trivial band-touching points that bring quantized numbers for integral of the Berry flux over any closed surface enclosing these points, and provide the relevant topological invariants. Depending on band degeneracy and the momentum-space distribution of the band-touching points, three different types of TSMs, namely, Weyl semimetal (WSM)<sup>5–10</sup>, Dirac semimetal (DSM)<sup>11,12</sup>, and nodal-line semimetal (NLSM)<sup>13–17</sup> have been proposed. WSMs and DSMs exhibit two- and four-fold band-touching points with low energy Weyl and Dirac fermion excitations, respectively, with zero dimensional (0D) Fermi surfaces in the 3D bulk Brillouin zone (BZ). In sharp contrast to WSMs and DSMs, NLSMs support extended band-touching points along a line with 1D Fermi surfaces in the bulk BZ. WSMs have been found in noncentrosymmetric TaAs<sup>9,10</sup>, Mo<sub>x</sub>W<sub>1–x</sub>Te<sub>2</sub><sup>18–20</sup>, LaAlGe<sup>21</sup> families, while DSMs have been realized in Na<sub>3</sub>Bi<sup>11</sup> and Cd<sub>2</sub>As<sub>3</sub><sup>12</sup>.

Among the TSMs, NLSMs offer many unique properties that are distinct from WSMs and DSMs<sup>14–17,22–24</sup>. For example, NLSMs have higher density of states at the Fermi level than DSMs and WSMs, and therefore, provide an ideal platform to study interaction-induced instabilities. They also feature topological surface states (TSSs), known as “Drumhead surface states (DSSs)”, which could be interesting for achieving superconductivity and correlation physics<sup>23–25</sup>. Despite much theoretical efforts<sup>26–34</sup>, the experimental evidence for NLSMs has been only reported recently for ZrSi(S,Te)<sup>35,36</sup>, PbTaSe<sub>2</sub><sup>37</sup>, and PtSn<sub>4</sub><sup>38</sup>. A focus of discussion has been the issue of the stability of nodal lines in the absence of spin-orbit coupling (SOC) effects<sup>16,17,26–29</sup>. Turning on the SOC either splits the nodal line into nodal points, depending on crystalline symmetries present<sup>7,8,26,27</sup>, or fully gaps the spectrum due to hybridization with bands of same symmetry<sup>28,29,34–36</sup>. Regardless, saddle-like TSSs have not been found in these materials.

In this paper, we identify topological nodal-line fermion states in the large family of silicides or germanides, TT’X (T, T’ = Transition metal, X= Si, or Ge) when the SOC is

ignored. Taking ZrPtGe as an explicit example, we show that a single nodal line lies on the  $k_y=0$  bulk plane and a saddle-like DSS nested inside the nodal line projection on the (010) surface. Inclusion of the SOC gaps the nodal line, realizing a  $Z_2$  nontrivial topological state with  $Z_2 = (1; 000)$ . The TSS evolves from the DSS with similar saddle-like energy dispersion. This is very unique since the known TIs such as  $\text{Bi}_2\text{Se}_3$  have Dirac-cone-like surface states<sup>39,40</sup>. It is well known that saddle-points in the band structure give rise to interesting saddle-point Van Hove singularities (VHSs) where the 2D density of states diverges<sup>41</sup>. When the VHS lies close to the Fermi level, the instabilities among lattice, charge, and spin degrees of freedom as well superconducting transition temperature, ferromagnetism and/or antiferromagnetism are substantially enhanced even in the weak coupling limit<sup>42–47</sup>. Our proposal of NLSM with saddle-like TSSs in the large family of silicides and germanides therefore provides an exciting materials platform to explore these exotic properties in the presence of a nontrivial band topology.

## II. METHOD AND CRYSTAL STRUCTURE

We performed electronic structure calculations with the projector augmented wave (PAW) method<sup>48–50</sup> and generalized gradient approximation (GGA)<sup>51</sup>, using the VASP<sup>49</sup>. SOC was included in calculations self-consistently. The surface energy dispersions were calculated within the tight-binding scheme, using the Wannier-tools<sup>52–54</sup>.

The TT'X compounds considered all crystallize in an orthorhombic Bravais lattice with the non-symmorphic space group  $D_{2h}^{16}$  ( $Pnma$ , No. 62)<sup>60,61</sup>. The crystal structure of ZrPtGe is illustrated in Fig. 1(a) as an example. In this structure, Pt and Ge atoms form a strongly corrugated  $\text{Pt}_3\text{Ge}_3$  hexagonal network and Zr atoms fill the cavities left in the network. Due to strong puckering between different atomic layers, Pt forms a distorted tetrahedral configuration with Ge whereas Zr is coordinated with five Ge atoms as seen in Fig. 1(b). The first orthorhombic bulk and (010) surface BZs with the relevant high-symmetry points are shown in Fig. 1(c).

### III. ELECTRONIC STRUCTURE

The electronic structure of ZrPtGe without the SOC [Fig. 1(d)] unveils its semimetallic ground state. The valence and conduction bands are seen to cross along the high-symmetry lines  $\Gamma - X$ ,  $\Gamma - Z$ , and  $\Gamma - U$  that are tied to the  $k_y = 0$  plane of the bulk BZ. A full BZ exploration shows that these band crossings persist along a closed path, realizing a single nodal line on the  $k_y = 0$  plane inside the bulk BZ as shown in Fig. 1(c). When we include the SOC, the nodal line evaporates with the opening of a gap at the band-crossing points [Fig. 1(e)]. It should be noted that owing to the coexistence of inversion ( $P$ ) and time reversal symmetry ( $\Theta$ ), each band still remains doubly degenerate in the presence of the SOC at each  $k$ -point. Furthermore, additional non-symmorphic crystalline symmetries in this system lead to four-fold band crossings at the BZ boundary planes ( $k_i = \pi$ ) above and below the Fermi level. These band crossings are protected against gap opening and may realize high-symmetry Dirac cones at  $X$ ,  $Y$ , and  $Z$  points. The energy-dispersion can be tuned by changing the transition metal elements in the system as shown in appendix.

In order to characterize the nature and the topological protection of the nodal line, we systematically examine the band crossings in Fig. 2. We know that nodal line resides on the  $k_y = 0$  plane, which is a  $M_y : (x, y, z) \rightarrow (x, -y, z)$  mirror invariant plane. Each band on this plane can have a well defined  $M_y$  mirror eigenvalue. If a band crossing happens between two bands of different eigenvalues, it can remain gapless. Figure 2(a) shows the bands along two principal directions on the  $k_y = 0$  plane and the corresponding mirror eigenvalues in the absence of the SOC. Note that we have obtained  $M_y$  eigenvalues from our first-principles Bloch wavefunctions and since the Hamiltonian remains spin-rotation invariant without SOC, it has eigenvalues  $+1$  or  $-1$ . Clearly, the lowest conduction band and the highest valence band have different mirror eigenvalues and thus their crossing points are topologically protected against gap opening. Further analysis of orbital characters shows that bands around the Fermi level are dominated by Pt and Zr  $d$  states with a clear signature of band inversion at the  $\Gamma$ -point [Fig. 2(a)]. We explicitly calculated the topological invariant for the system using the mirror eigenvalue analysis<sup>28,29</sup>. The computed topological invariant  $\nu$  takes nontrivial value only inside the nodal line (Fig. 2(c)) and thus, signals the existence of an odd number of DSSs inside the nodal line projection over the surface<sup>28,29</sup>.

The topological stability can be further assessed by calculating the topological invariants

$Z_2 = (\nu_0; \nu_1 \nu_2 \nu_3)$  defined in the presence of  $P$  and  $\Theta$  symmetries without the SOC<sup>26</sup>. Our calculations yield a nontrivial  $Z_2 = (1;000)$  [see Fig. 2(d)]. This result demonstrates that the nodal-line in ZrPtGe is also protected by  $P$  and  $\Theta$ . In Fig. 2(b), we present the nodal line structure in the  $E - k_x - k_z$  space. The nodal line is seen to exhibit a substantial energy dispersion around the Fermi level with the corresponding energies of 0.189 eV, -0.023 eV, and -0.119 eV along  $k_x$ ,  $k_z$ , and the plane diagonal to the  $(k_x - k_z)$  directions. The large energy dispersion of the nodal line results in a unique Fermi surface that constitutes compensated electron and hole pockets as shown in Fig. 2(e).

Figure 2(f) illustrates the energy bands with SOC. In general SOC can drive NLSMs into DSM, WSM, or a fully gapped insulator. However, in view of the lowered crystalline symmetries of ZrPtGe, the SOC opens a full gap at the crossing points, making the conduction and valence band to separate at each k-point. Since the nodal line winds around a single time reversal invariant point ( $\Gamma$ ) with a nontrivial band topology, the SOC drives the system into a nontrivial insulating state with  $Z_2=(1;000)$  [see appendix for details].

To further showcase the protected surface states and their connection to bulk nodal-lines, we present these states for the semi-infinite (010) surface in Fig. 3. The bulk bands projected onto the (010) surface without the SOC are shown in Fig. 3(a) where the nodal-line crossings can be clearly seen. The DSSs nested inside the nodal line are visible in Fig. 3(b), which is consistent with the calculated non-trivial invariant inside the nodal line. Unlike the nearly flat DSSs reported in earlier works<sup>26-29</sup>, the states in ZrPtGe are more dispersive, and interestingly, they have opposite band curvatures along the  $\bar{\Gamma} - \bar{X}$  and  $\bar{\Gamma} - \bar{Z}$  directions, realizing a unique saddle-like energy dispersion.

Figures 3(c) and 3(d) show (010)-projected bulk bands and surface bands, respectively, with the SOC. The nodal line is now gapped and the DSS splits away from the time-reversal invariant  $\bar{\Gamma}$ -point, deforming into a topological Dirac-cone state. Since the SOC here is much smaller than the dispersion of the DSS, the upper and lower branches of the TSS display the same band curvatures or carrier velocities. Furthermore, these states retain the saddle-like features of the DSS energy dispersion, which is more clearly visible in the Fermi band contours shown in Figs. 3(e)-3(f). The constant energy contours (CECs) are seen to be open and disperse along the  $k_x$  direction above the Dirac-point or the saddle-point. As we lower the Fermi energy, the electronic states undergo a Lifshitz transition and the CECs change direction and dispersion to lie along the  $k_z$  direction.

We want to emphasize that TSSs with saddle-like energy dispersion are unique to the ZrPtGe family. They are symmetry allowed (see below) and their characteristic saddle-like dispersion is controlled by their bulk nodal-line structure. Therefore, we should expect these states to be robust over a range of different surface potentials, although details of their energy positions can be shifted via changes in surface potentials, see appendix for details. In particular, we have found that when we repeat our computations using the more advanced Heyd-Scuseria-Ernzerof (HSE) exchange-correlation functional to obtain the bulk and surface spectrum, the dispersive nature of the bulk nodal-line as well as the saddle-like TSSs remain preserved.

#### IV. LOW-ENERGY EFFECTIVE MODEL

In order to better understand the saddle-like TSSs, we now present a low-energy effective model Hamiltonian using the theory of invariants in a manner similar to the case of Bi<sub>2</sub>Te<sub>3</sub><sup>40</sup>. On the (010) surface, in addition to  $\Theta$ , the only other preserved symmetry is the glide-mirror symmetry  $\bar{M}_z = \{M_z | \frac{1}{2}0\frac{1}{2}\}$ , which sends  $(k_x, k_z) \rightarrow (k_x, -k_z)$ . Based on our first-principles results and a symmetry analysis, a single-band  $k \cdot p$  model Hamiltonian for the DSSs in the absence of the SOC takes the form

$$H_0(k_x, k_z) = \frac{1}{2m^*}(k_x^2 - \eta k_z^2) = -\frac{1}{2m^*}[\frac{\eta+1}{4}(k_+^2 + k_-^2) + \frac{\eta-1}{2}k_+k_-], \quad (1)$$

where  $k_{\pm} = k_z \pm ik_x$ . Here coefficient  $\eta$  describes the form of the  $E - K$  dispersion and depends on the rotational symmetries and materials properties of the system. While  $\eta < 0$  gives a parabolic energy dispersion,  $\eta > 0$  ensures a saddle-like energy dispersion (see appendix for details). The  $n$ -fold rotational symmetry  $C_n$  for  $n > 2$  normal to the surface in a system forbids  $\eta > 0$  since it requires  $\eta = -1$ , i.e a saddle-like energy dispersion for the surface state whereas for  $n \leq 2$  this is allowed. We have verified this point from symmetry constrains of  $C_n$  on the Hamiltonian  $H_0(k_x, k_z)$  with  $\eta > 0$ . Since  $C_n$  sends  $k_{\pm} \rightarrow k_{\pm} e^{\pm i \frac{2\pi}{n}}$  and  $(k_+^2 + k_-^2) \rightarrow (k_+^2 e^{i \frac{4\pi}{n}} + k_-^2 e^{-i \frac{4\pi}{n}})$ , this Hamiltonian remains invariant only for  $C_n$  with  $n \leq 2$ . Notably, this is only a necessary condition for realizing a saddle-like energy dispersion as material properties are also involved in achieving such a dispersion. As ZrPtGe exhibits a substantial nodal-line dispersion in the bulk and its (010) surface lacks  $C_{ny}$  with  $n > 2$ , it hosts a symmetry allowed saddle-like state as found in our results [cf. Fig. 3]. The energy

dispersion associated with the model Hamiltonian 1 is presented in Figs. 4(a)-4(c), which show a saddle-like energy dispersion for the DSS with a single saddle-point at the  $\Gamma$ -point. The density of states (DOS) is logarithmically diverging at the  $\Gamma$  point, confirming that it has a saddle-point VHS.

In the presence of SOC, the DSS splits into two branches, developing into the spin-polarized surface states of a strong  $Z_2$  TI. A two-band  $k.p$  Hamiltonian is therefore necessary for describing these states. Considering  $\overline{M}_z = -e^{-i(k_x - k_z)/2} i\sigma_z$  and  $\Theta = i\sigma_y K$ , where  $\sigma_{i=x,y,z}$  denote Pauli spin matrices, and  $K$  is the complex conjugate operator, the effective model Hamiltonian with SOC can be written as

$$H_{SOC}(k_x, k_z) = \frac{1}{2m^*}(k_x^2 - \eta k_z^2) + v_k(k_z\sigma_x - k_x\sigma_z) - v'_k k_z\sigma_y, \quad (2)$$

where  $v_x = v_k = v_0(1 + \alpha k^2)$  and  $v_z = \sqrt{v_k^2 + v'_k{}^2} = v_{z0}(1 + \alpha k^2)$  are the Dirac velocities along the  $x$ - and  $z$ -axis, respectively, with a second order correction. The corresponding energy dispersion  $E_{\pm}(\mathbf{k})$  is

$$E_{\pm}(\mathbf{k}) = \frac{1}{2m^*}(k_x^2 - \eta k_z^2) \pm \sqrt{v_x^2 k_x^2 + v_z^2 k_z^2}. \quad (3)$$

The preceding energy dispersion demonstrates a new type of symmetry-allowed TSS which is distinct from surface states studied so far<sup>1,2</sup>. This new TSS evolves from a saddle-like DSS with a single Dirac-point at  $\Gamma$  and two pairs of saddle points at generic  $k$  points. The two saddle points are located at  $(k_x, k_z) = (\pm m^* v_x, 0)$  with energy,  $\omega_-^* = -\frac{m^* v_x^2}{2}$ , on the lower branch while the other two points lie at  $(k_x, k_z) = (0, \pm m^* v_z/\eta)$  with energy  $\omega_+^* = \frac{m^* v_z^2}{2\eta}$  on the higher branch of the TSS (see Figs. 4(d)-4(f)). This can be further seen in the DOS [Fig. 4(f)] where the two saddle-point VHSs are evident at  $\omega_-^*$  and  $\omega_+^*$ . Such VHSs which appear because of the saddle-points at generic  $k$ -points are classified as type-II VHSs<sup>46</sup>.

## V. DISCUSSION

The exotic TSSs with VHSs that we have delineated in this study would provide a new platform for exploring the interplay between topological states and strong correlation physics and the related interaction-driven instabilities. For example<sup>46</sup>, a type-II VHS system can favor the odd-parity pairing and lead to unconventional superconductivity via weak repulsive interactions. Also, the generic competing orders such as the "valley" charge imbalance,



density wave orders (through enhanced nesting via a uniaxial pressure, for example), and superconductivity could possibly be driven by electron-electron interactions<sup>23–25,42–47</sup>. It should be emphasized, however, that the actual ground state realized by a specific material will be influenced by details of the location of the chemical potential, nesting properties, and the nature of the various interactions involved and their relative strengths.

The transition metal silicides and germanides have been explored in connection with search for new high-temperature superconductors in the intermetallics. In particular, a number of the materials we have proposed here have been synthesized as single crystals<sup>57,58,60,61</sup>. Since our predicted exotic states lie close to the Fermi level, these states could of course be probed directly via ARPES experiments. Moreover, the proposed materials feature unique bulk Fermi surfaces, which may, for example, lead to balanced electron-hole resonance conditions, and thus induce unusual transport characteristics such as a large positive unsaturated magnetoresistance<sup>20</sup>.

In summary, our first-principles computations predict that the orthorhombic TT'X family of silicides and germanides harbors a single topologically protected nodal line on the  $k_y = 0$  plane in the absence of the SOC. We have demonstrated the existence of DSSs nested inside the nodal line on the (010) surface of ZrPtGe as an exemplar system, which hosts a unique symmetry-allowed saddle-like energy-momentum dispersion relation. Inclusion of the SOC gaps the nodal line and eventually drives the material into the TI phase. The nontrivial TSSs evolve from the DSS and retain their saddle-like energy dispersion and support two pairs of saddle point VHSs. Our results establish that TT'X materials family provides an ideal platform for exploring unique physics, including symmetry-breaking quantum states, related with the NLSMs and the saddle-like topological surface states.

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## APPENDIX A: BAND-GAP CORRECTION

It is well known that the generalized gradient approximation (GGA) usually underestimates band gaps and results in an overestimation of the band inversion strength. In order to check the robustness of our band inversion and the nodal-line semimetal state in ZrPtGe, we also carried out computations using the more advanced Heyd-Scuseria-Ernzerof (HSE) exchange-correlation functional<sup>59</sup>, which yields improved band gaps in closer agreement with experiments. Fig. 5(a) shows the band structure of ZrPtGe obtained using the GGA, which is the band structure presented also in the main text. The nodal-line band crossings are highlighted with broken red circles. Fig. 5(b) gives the band structure based on the HSE functional, where although the band gaps at high-symmetry points are slightly enlarged, the overall features of the GGA band structure are seen to be retained. These results demonstrate that the topological properties and nodal-line semimetal state of ZrPtGe is quite robust to changes in the exchange-correlation functional.

In Figs. 5(c) and 5(d), we present the topological surface states of ZrPtGe on the (010) surface obtained with GGA and HSE functional, respectively. The non-trivial drumhead surface states (DSSs) with unique saddle-like energy dispersion shown in Fig. 5(c) are same as those presented in the main text. Similar to the bulk nodal-line structure, the DSSs obtained with HSE continue to display saddle-like energy dispersion, see 5(d).

## APPENDIX B: NON-TRIVIAL BAND TOPOLOGY

Here we present calculations of non-trivial topological invariants of ZrPtGe in the presence of spin-orbit coupling (SOC). As discussed in the main text, ZrPtGe is a non-trivial semimetal with  $Z_2 = (1;000)$  without the SOC. While the semimetal state is maintained when the SOC is included with the presence of electron and hole pockets at the Fermi level, the SOC separates valence and conduction bands at all crossings points and opens a contin-

uous band gap between valence and conduction states, see Figs. 1 and 2 of the main text. The  $Z_2 = (\nu_0; \nu_1\nu_2\nu_3)$  topological invariants can still be determined as in a fully gapped insulator. Fig. 6 shows the evolution of Wannier charge centers (WCCs) in the six time-reversal invariant momentum planes in the bulk Brillouin zone (BZ) of ZrPtGe. The WCCs have non-trivial connectivity in  $k_i = 0$  planes with  $Z_2$  index 1, while they are connected trivially in  $k_i = \pi$  planes with  $Z_2 = 0$ , leading to  $Z_2=(1;000)$  or a strong topological insulator phase in ZrPtGe.

### APPENDIX C: ROBUSTNESS OF SADDLE-LIKE TOPOLOGICAL SURFACE STATES

ZrPtGe has a layered crystal structure along the (010) direction with two repeated atomic layers, see Fig. 7. Each atomic layer contains all three constituent atoms (Zr, Pt, and Ge). These two atomic layers are related via  $\{C_{2y}|0\frac{1}{2}0\}$ . Unlike TlBiSe<sub>2</sub><sup>62</sup>, if we consider only a flat surface, there is only one surface termination, which is highlighted by the broken blue line in Fig. 7. This is the surface termination we use in our calculations.

We have examined the robustness of topological surface states and their saddle-like dispersion in ZrPtGe within the framework of our tight-binding Hamiltonians by including an additional surface potential energy (SPE) as an on-site energy parameter, which is added to all the atoms in the surface layer. Figs. 8 and 9 present the results obtained by changing the SPE from -0.3 eV to +0.3 eV for our GGA-based tight-binding Hamiltonian. Although the surface states shift towards the bulk valence or conduction bands depending on the sign of the SPE parameter, we see that the saddle-like dispersion of the surface states remains unchanged. These saddle-like surface states can be removed if they hybridize with the bulk bands, which occurs for large values of the SPE ( $\sim \pm 0.4$  eV), much like the Dirac surface states in Bi<sub>2</sub>Se<sub>3</sub>. Similar results are obtained when we consider effects of the value of the SPE parameter on our HSE-based tight-binding Hamiltonian. These results show clearly that the TSSs with saddle-like dispersion in ZrPtGe are robust against variations in the surface potential.

## APPENDIX D: MODEL HAMILTONIAN FOR THE (010) SURFACE STATES

The effective  $k, p$  model Hamiltonian for the saddle-like TSSs in the presence of the SOC and the corresponding energy dispersion given in Eqs. (2) and (3) of the main text are:

$$H_{SOC}(k_x, k_z) = \frac{1}{2m^*}(k_x^2 - \eta k_z^2) + v_k(k_z\sigma_x - k_x\sigma_z) - v'_k k_z\sigma_y, \quad (4)$$

$$E_{\pm}(\mathbf{k}) = \frac{1}{2m^*}(k_x^2 - \eta k_z^2) \pm \sqrt{v_x^2 k_x^2 + v_z^2 k_z^2}. \quad (5)$$

Here,  $v_x = v_k = v_0(1 + \alpha k^2)$  and  $v_z = \sqrt{v_k^2 + v'_k{}^2} = v_{z0}(1 + \alpha k^2)$  are the Dirac velocities along the  $x$ - and  $z$ -axis, respectively, including a second order correction. The first and second terms in Eq. 4 describe the kinetic energy and the Rashba SOC, respectively. The third term arises from the lowered symmetry group of the (010) surface. Its origin can be understood by recalling how the SOC comes about. When an electron with momentum  $\mathbf{k}$  moves in an electric field  $\mathbf{E}$ , it experiences a magnetic field  $\mathbf{B}_{eff} \sim \mathbf{E} \times \mathbf{k}/mc^2$  in its rest frame. This gives rise to a momentum-dependent Zeeman interaction or the SOC,  $\hat{H}_{SO} \sim \mu_B(\mathbf{E} \times \mathbf{k}) \cdot \boldsymbol{\sigma}/mc^2$ . In a crystal, the electric field is given by the gradient of the crystal potential  $\mathbf{E} = -\nabla V$ <sup>65</sup>.

For a 2D system with  $C_{2v}$  symmetry, Vasko<sup>63</sup> and Bychkov and Rashba<sup>64</sup> showed that the interfacial electric field perpendicular to the surface,  $\mathbf{E} = E_z \hat{\mathbf{z}}$ , gives rise to the SOC of the form  $\hat{H}_R = \frac{\alpha_R}{\hbar}(\hat{\mathbf{z}} \times \mathbf{k}) \cdot \boldsymbol{\sigma} = \frac{\alpha_R}{\hbar}(k_x\sigma_y - k_y\sigma_x)$ , where  $\alpha_R$  is called the Rashba parameter. In ZrPtGe, the bulk system respects a two-fold screw rotation symmetry,  $\{C_{2y}|0\frac{1}{2}0\}$ , which yields two symmetry related surface terminations on the (010) surface. However, for either surface termination, the symmetry is lowered to  $C_s$  with only one glide mirror plane  $\{M_z|\frac{1}{2}0\frac{1}{2}\}$ . This allows an interfacial electric field  $\mathbf{E} = (E_x, E_y, 0)$  and results in the survival of the third term in Eq. 4 on the surface. It should be noted that for the other surface termination, the sign of the third term in Eq. 4 is reversed.

We now discuss the classification of saddle-points and possible Lifshitz transitions resulting from variations in  $\eta$  based on our model Hamiltonian. As noted above as well as in the main text, the first-principles (010) surface of ZrPtGe supports the presence of unique surface states with a saddle-like dispersion, which is described by  $\eta > 0$  in our model Hamiltonian. Without SOC,  $\eta > 0$  guarantees a saddle-point Van Hove singularity (VHS) at  $\Gamma$ . Moreover, when the SOC is included,  $\eta > 0$  induces two pairs of type-II VHSs on the  $k_x$  and  $k_z$  axis. Note that VHSs in ZrPtGe can only appear on the  $k_x$  or  $k_z$  axis due to the

constraints of  $C_s$  and time-reversal symmetries.

In Fig. 10, we demonstrate a general classification of the saddle-point VHSs as a function of  $\eta$  for a 2D system with  $C_n$  for  $n \leq 2$ , taking linear SOC terms into account.  $\eta = -v_z^2/v_x^2$  gives a typical Rashba-like parabolic energy dispersion without a saddle-point VHS. However, for  $\eta > -v_z^2/v_x^2$ , one pair of saddle-point VHSs appears on the  $x$ -axis at  $(k_x, k_z) = (\pm m^* v_x, 0)$  with energy  $\omega_1^* = -\frac{m^* v_x^2}{2}$ . For  $\eta < -v_z^2/v_x^2$  or  $\eta > 0$ , one pair of VHSs emerges on the  $z$ -axis at  $(k_x, k_z) = (0, \pm m^* v_z/\eta)$  with energy  $\omega_2^* = \frac{m^* v_z^2}{2\eta}$ . In other words, if  $\eta \leq 0$  (excepting  $\eta = -v_z^2/v_x^2$ ), one pair of saddle-point VHSs can still be induced by the SOC. Note, however, that this is a consequence of the SOC, so that no saddle-point VHSs can appear in the absence of the SOC for  $\eta \leq 0$ . Finally, for  $\eta > 0$ , two pairs of saddle-point VHSs emerge with one pair lying on the  $x$ -axis and the other on the  $z$ -axis, as is the case in ZrPtGe.

For a 2D system ( $xz$ -plane) with  $n$ -fold rotation symmetry about the  $y$ -axis,  $C_{ny}$ , for  $n \geq 2$ , the allowed linear SOC terms are  $i^\lambda v_\lambda (k_+ \sigma_- + (-1)^\lambda k_- \sigma_+)$ , where  $\lambda = 0$  or  $1$ . Here,  $C_{ny}$  sends  $\sigma_\pm \rightarrow e^{\pm i \frac{2\pi}{n}} \sigma_\pm$  and  $k_\pm \rightarrow e^{\pm i \frac{2\pi}{n}} k_\pm$ , where  $k_\pm = k_z \pm i k_x$  and  $\sigma_\pm = \sigma_z \pm i \sigma_x$ . The presence of additional crystalline symmetries drives the SOC term into taking the form of the Rashba SOC,  $v_1(k_z \sigma_x - k_x \sigma_z)$  for  $\lambda = 1$  ( $C_{nv}$  symmetry, for example);  $v_0(k_z \sigma_z + k_x \sigma_x)$  for  $\lambda = 0$  (e.g.  $D_n$  symmetry); and, both  $v_0$  and  $v_1$  terms are allowed (e.g.  $C_n$  symmetry). Note that according to Eq. 5, in all cases  $v_z^2 = v_x^2$  ( $= v_1^2$ , or  $v_0^2$ , or  $v_0^2 + v_1^2$ ), resulting in the critical point  $-v_z^2/v_x^2 = -1$ . For  $n > 2$ , however,  $\eta$  is not allowed to take any other value than  $-1$ . This clearly shows that a system with  $C_{ny}$  ( $n > 2$ ) can only support Rashba-like parabolic energy dispersion for the surface state when only linear SOC terms are considered.

## APPENDIX E: BAND STRUCTURE OF TT'X FAMILY

Figs. 11 and 12 present the bulk band structures of all members of the TT'X family without and with SOC. The 16 members in this family are listed in Table I along with the lattice constants<sup>60,61</sup> used in our calculations and their topological invariants (with SOC). All these compounds host single nodal lines (without SOC) with substantial dispersions similar to ZrPtGe. Inclusion of SOC gaps the nodal line in all cases and yields the topological insulator state with  $Z_2=(1;000)$ . A reference to the band structures, however, shows that the area occupied by the nodal loop on the  $k_y = 0$  plane of the bulk BZ depends on the transition

metal element involved. This is also the case with the details of band crossings at the high-symmetry points above and below the Fermi level (see Figs. 11 and 12). These results show that the TT'X family offers an interesting materials platform where the electronic structure can be tuned through the choice of the transition metal element.

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\* [bahadursingh24@gmail.com](mailto:bahadursingh24@gmail.com)

† [nilnish@gmail.com](mailto:nilnish@gmail.com)

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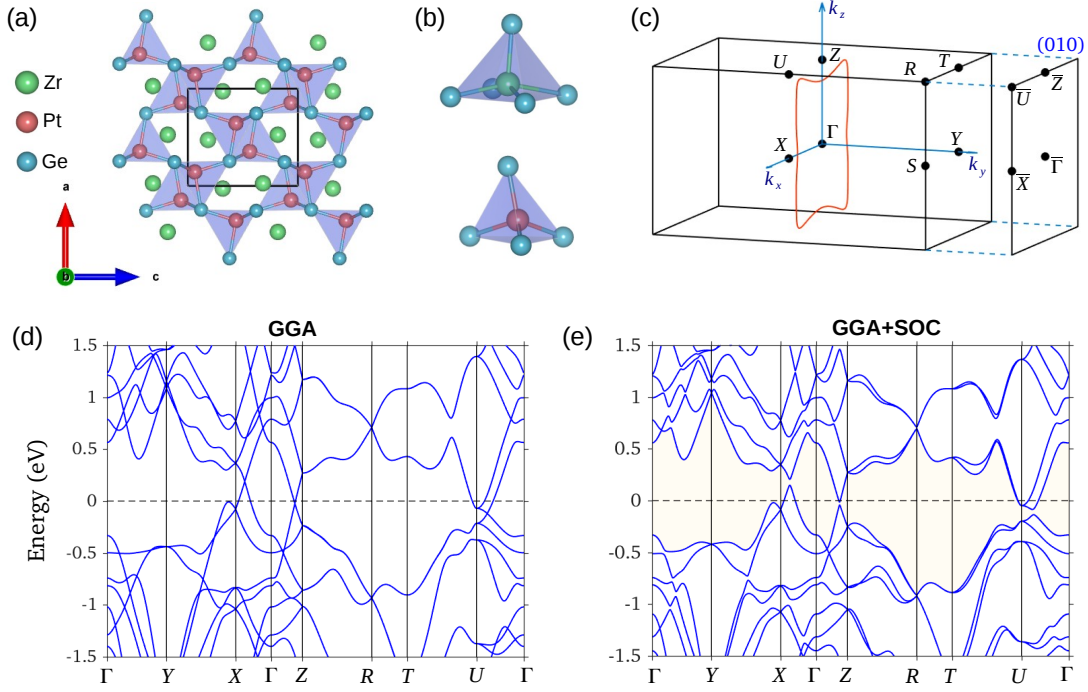


FIG. 1. (a) Crystal structure of ZrPtGe with  $Pnma$  symmetry. (b) Local coordination of the Zr (top panel) and Pt (bottom panel) atoms in the unit cell. (c) Bulk Brillouin zone and its projection on the (010) surface. The relevant high-symmetry points are marked with black dots. The calculated momentum distribution of the nodal line is shown with orange on the  $k_y=0$  plane. (d) Bulk band structure without spin-orbit coupling (SOC). Band crossings are evident along the  $\Gamma - X$ ,  $\Gamma - Z$ , and  $\Gamma - U$  directions near the Fermi level. (e) Bulk band structure with SOC. A small band gap opens at the band-crossing points. The shaded region highlights the continuous gap below which the topological invariants are calculated.

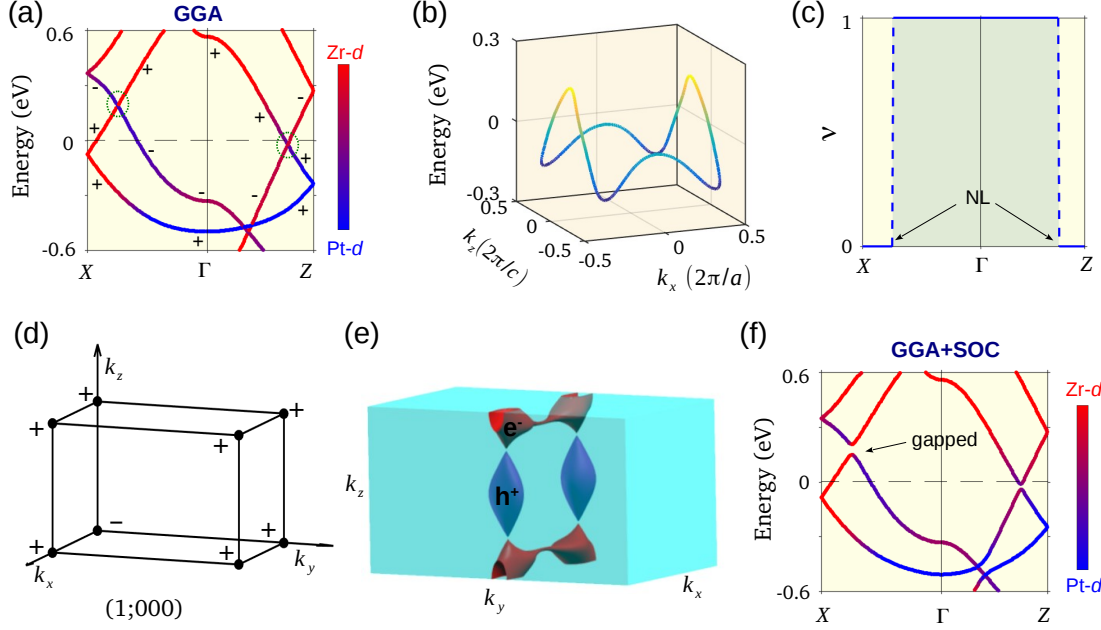


FIG. 2. (a) Orbital compositions and mirror eigenvalues of the bulk bands without SOC. (b) Energy-dependent nodal line configuration on the  $k_y=0$  plane. (c) Variation of non-trivial topological invariant along the high-symmetry lines in the bulk Brillouin zone. The topological invariant is calculated using  $M_y$  mirror eigenvalues of the occupied bands on the mirror symmetric  $k_y = 0$  and  $k_y = \pi$  planes without SOC. (d) Products of parity eigenvalues of the occupied bands at eight time-reversal-invariant  $k$  points in the bulk BZ and the related  $Z_2$  indices without SOC. (e) The Fermi surface with electron (red) and hole (violet) pockets. (f) Orbital compositions of the bulk bands with SOC. A small band gap opens at the band crossings.

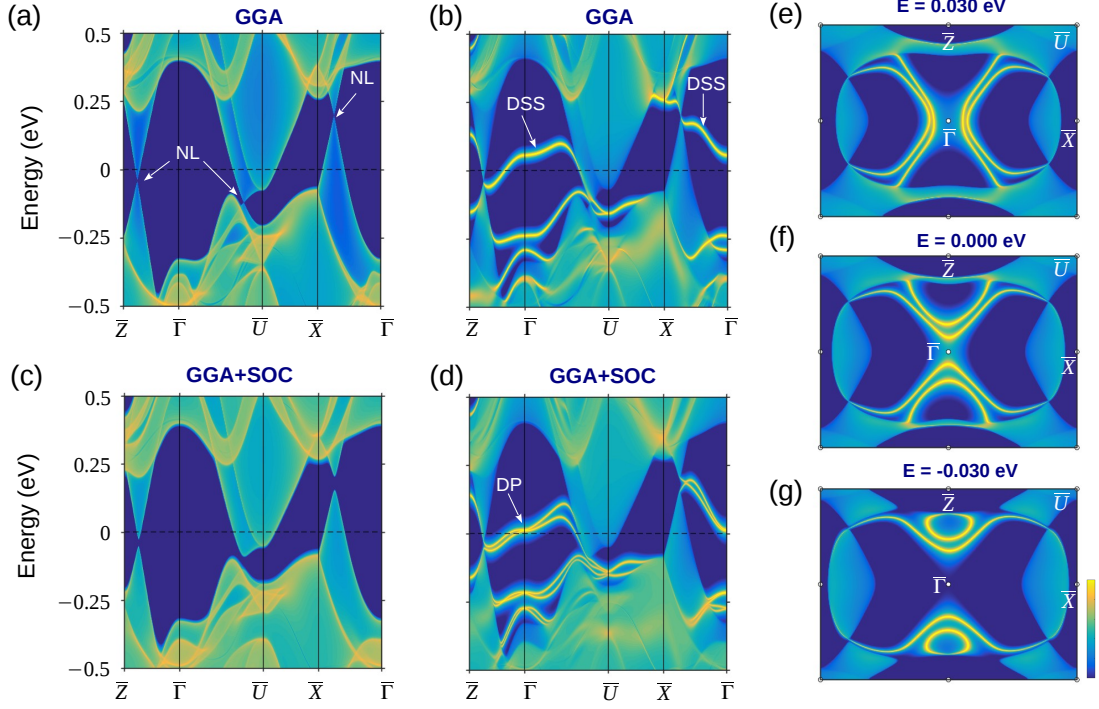


FIG. 3. (a) Projected bulk bands onto the (010) surface without SOC. The nodal line (NL) crossings are shown with arrows. (b) Band structure of the (010) surface without SOC. The non-trivial drumhead surface states (DSSs) are marked with arrows. (c)-(d) Same as (a)-(b) but with the SOC included. The non-trivial surface states evolve from the DSSs with different band curvatures along different high-symmetry directions. The Dirac point (DP) is identified with the arrow. Isoenergy band contours at (f)  $E = 0.030$  eV, (g)  $E = 0.000$  eV, and (h)  $E = -0.030$  meV with SOC.

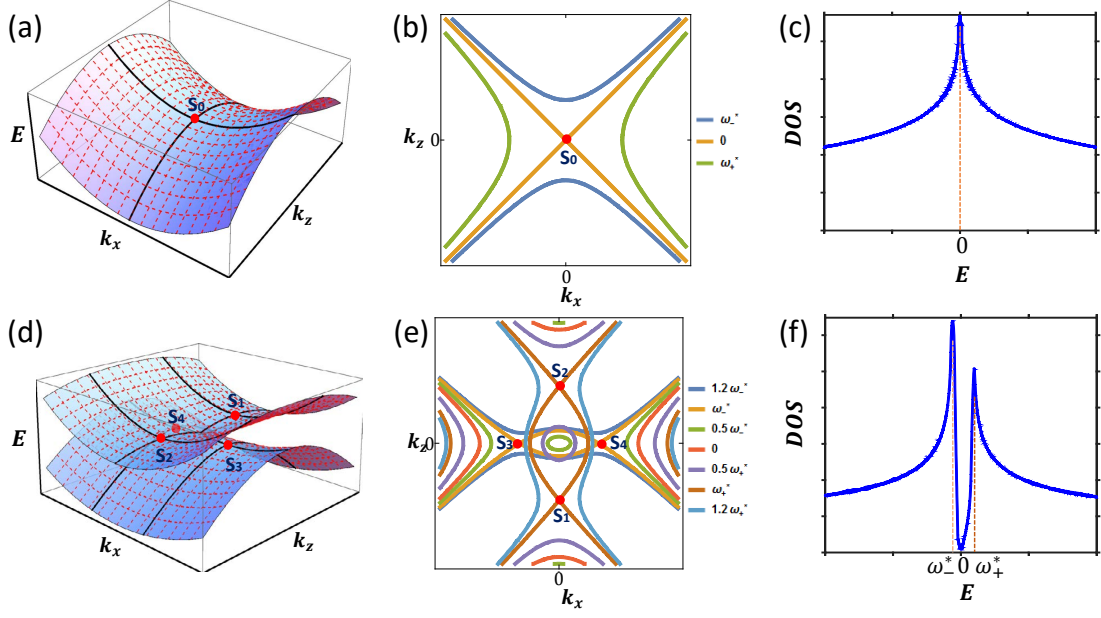


FIG. 4. (a) Energy dispersion of the Hamiltonian  $H_0(k_x, k_z)$  with  $\eta = +1$ , and (b) the associated isoenergy band contours without SOC. A single saddle-point  $S_0$  of the DSS lies at the  $\Gamma$  point. See main text for the meaning of energy band contours  $\omega_+^*$  and  $\omega_-^*$ . (c) Calculated density of states (DOS) for the surface states using the model Hamiltonian (eq. 1). (d)-(e) Same as (a)-(b) but with the SOC included. Four saddle points away from  $\Gamma$  ( $S_1, S_2, S_3$ , and  $S_4$ ) emerge due to band splittings in the presence of the SOC. (f) DOS for the surface states with SOC (eq. 4).

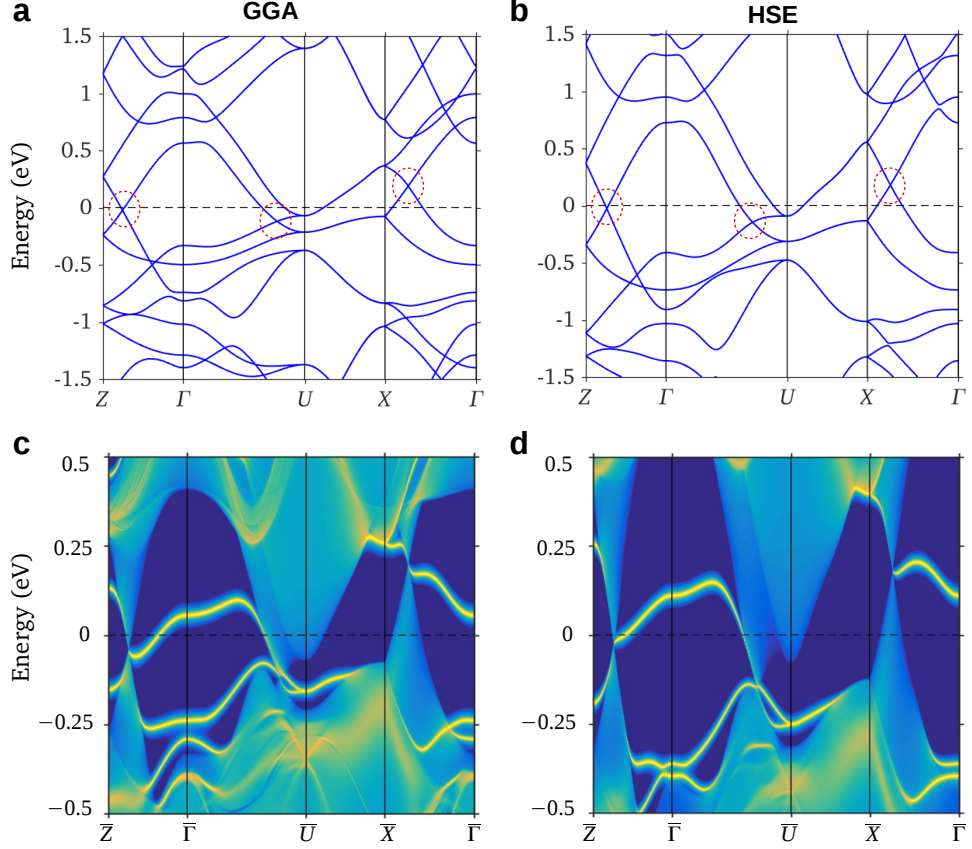


FIG. 5. Band structure of ZrPtGe without SOC. (a) The GGA band structure showing the nodal-line semimetal state. (b) The HSE band structure. The nodal-line crossings along the high-symmetry lines are highlighted with broken-red circles and remain intact after the band gap correction. Band structure of the (010) surface obtained using the (c) GGA and (d) the HSE functional. The non-trivial drumhead surface states (DSSs) and their saddle-like energy dispersion maintained with the HSE functional.



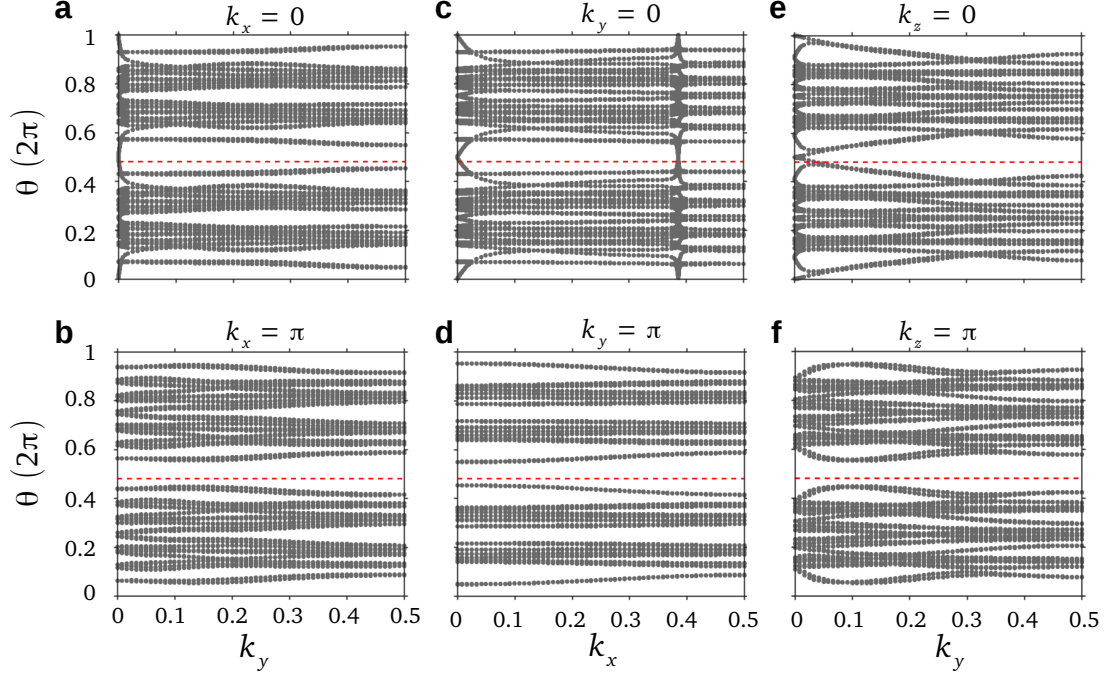


FIG. 6.  $Z_2$  topological invariants for ZrPtGe. Evolution of Wannier charge centers (WCCs) for (a)  $k_x = 0$ , (b)  $k_x = \pi$ , (c)  $k_y = 0$ , (d)  $k_y = \pi$ , (e)  $k_z = 0$ , and (f)  $k_z = \pi$  time-reversal invariant momentum planes. WCC evolution lines cross an arbitrary reference line (red broken line) an odd number of times on the  $k_i = 0$  plane, resulting in  $Z_2 = 1$ .

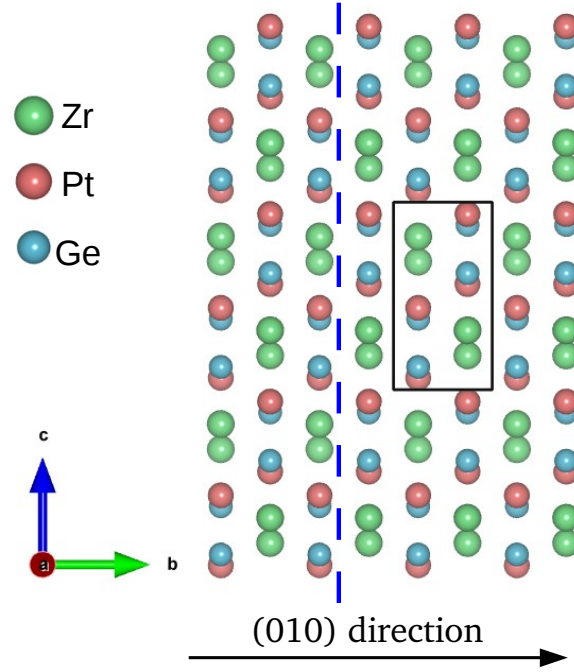


FIG. 7. Single surface termination along the (010) direction in ZrPtGe, which is marked with the dashed blue line. Black rectangle encloses the bulk unit cell in the  $yz$ -plane.

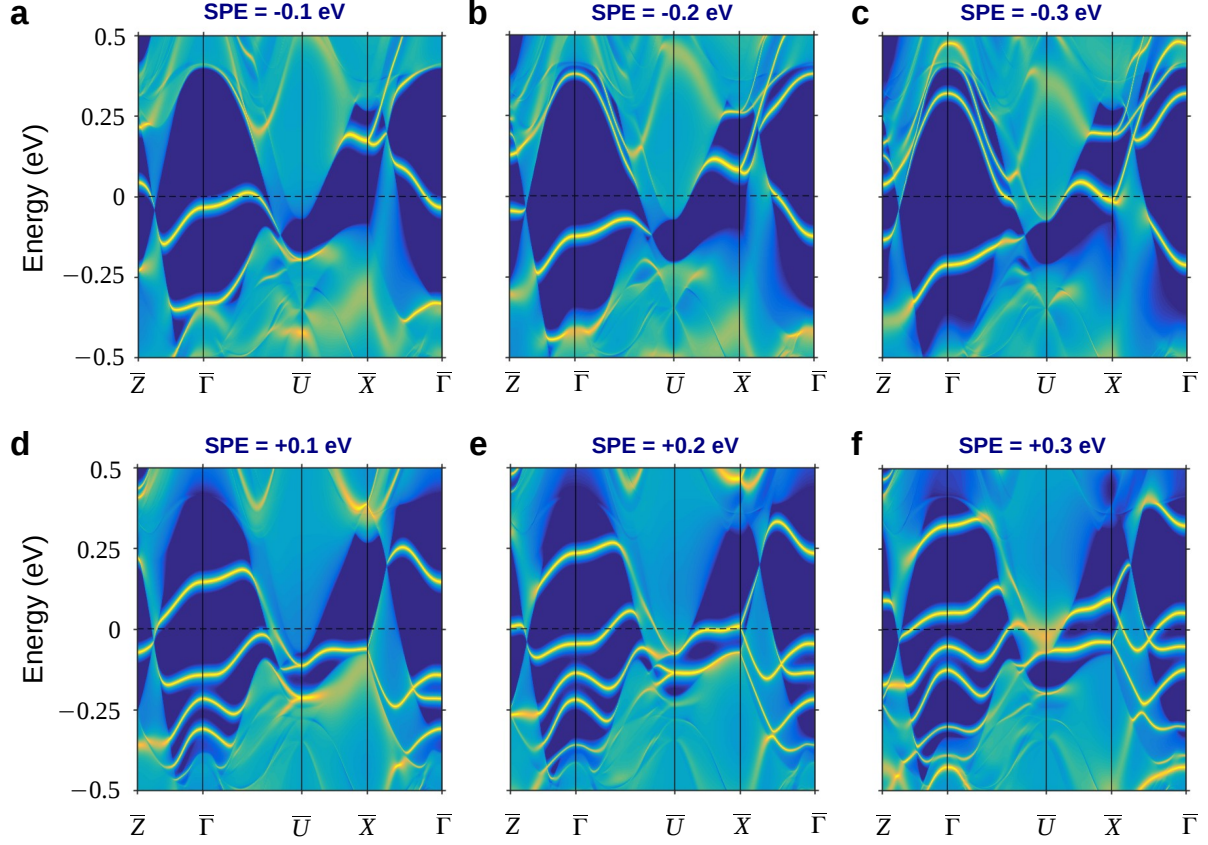


FIG. 8. Surface states of ZrPtGe (without SOC) obtained with different values of the surface potential energy (SPE) parameter: (a) -0.1 eV, (b) -0.2 eV, (c) -0.3 eV, (d) +0.1 eV, (e) +0.2 eV, and (f) +0.3 eV.

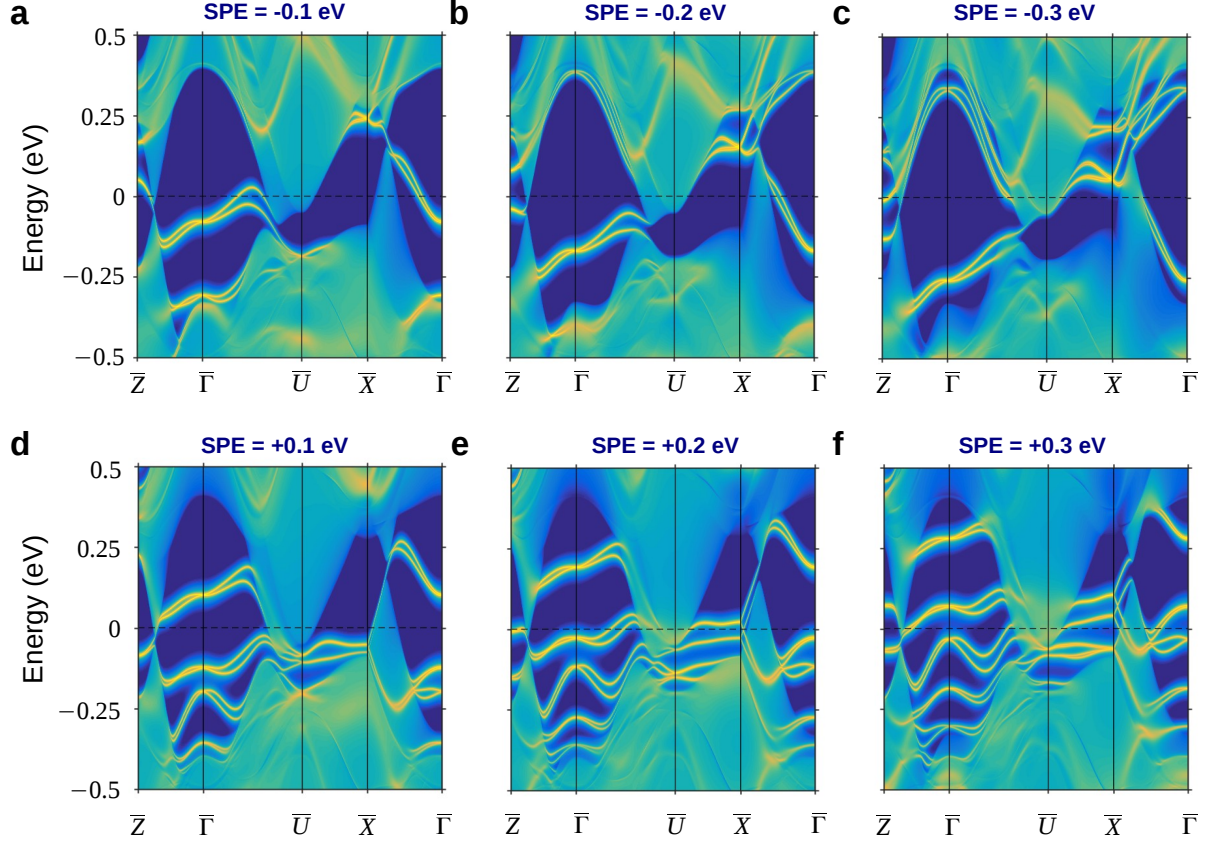


FIG. 9. Same as the caption of Fig. 8, except that this figure refers to results in which the SOC is included.

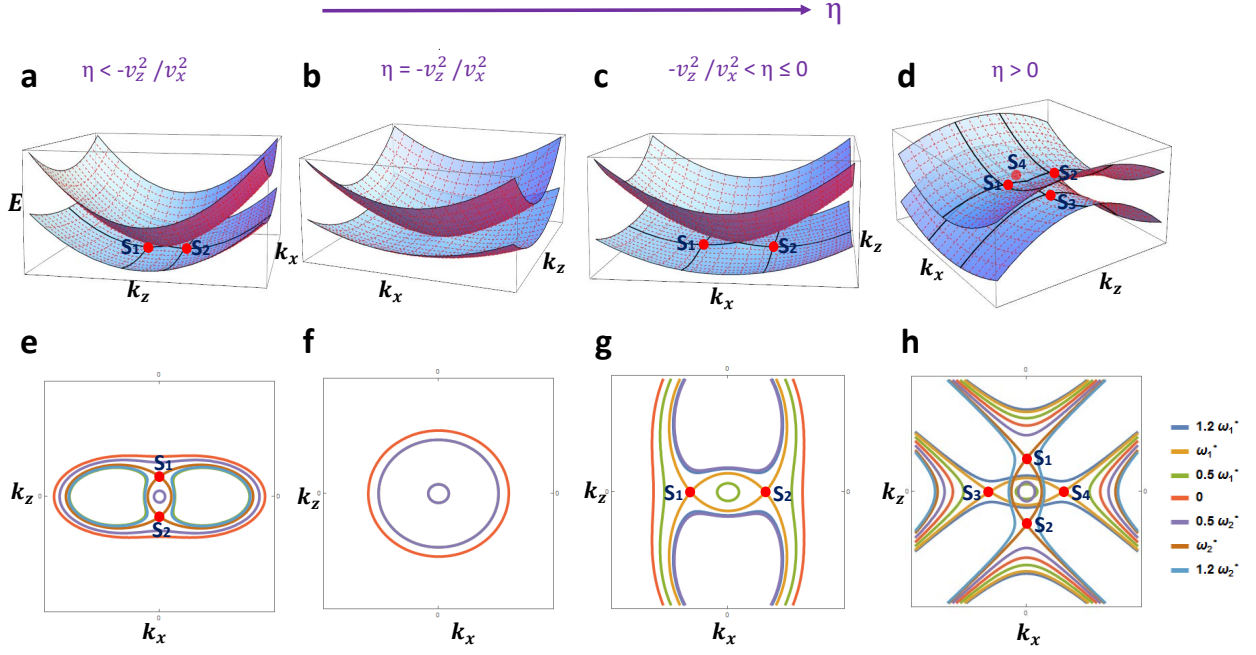


FIG. 10. Transitions in surface states as a function of the parameter  $\eta$ . (a-d) Energy dispersions, and (e-h) the associated isoenergy band contours for the Hamiltonian  $H_{soc}(k_x, k_z)$ ;  $\eta$  values are given in the top row. VHSs ( $S_1$ ,  $S_2$ ,  $S_3$ , and  $S_4$ , shown by red dots) emerge in pairs on the  $k_x$  or  $k_z$  axis for  $\eta \neq -v_z^2/v_x^2$  (see text for details).

TABLE I. Lattice constants  $a$ ,  $b$ , and  $c$  used in our calculations and the  $Z_2$  invariants for all 16 members of the TT'X family of compounds. The lattice parameters are taken from experimental studies<sup>60,61</sup>.

Compound	$a$ (Å)	$b$ (Å)	$c$ (Å)	$Z_2$
ZrPtGe	6.658	3.975	7.665	(1;000)
ZrPdGe	6.677	3.954	7.702	(1;000)
ZrNiGe	6.553	3.883	7.360	(1;000)
ZrPtSi	6.597	3.902	7.539	(1;000)
ZrPdSi	6.590	3.890	7.570	(1;000)
ZrNiSi	6.470	3.815	7.263	(1;000)
TiPtGe	6.371	3.857	7.537	(1;000)
TiNiGe	6.244	3.747	7.147	(1;000)
TiPtSi	6.336	3.802	7.337	(1;000)
TiPdSi	6.324	3.779	7.394	(1;000)
TiNiSi	6.139	3.661	7.006	(1;000)
HfPtGe	6.603	3.950	7.617	(1;000)
HfNiGe	6.500	3.810	7.290	(1;000)
HfPtSi	6.550	3.881	7.505	(1;000)
HfPdSi	6.570	3.874	7.675	(1;000)
HfNiSi	6.390	3.890	7.200	(1;000)



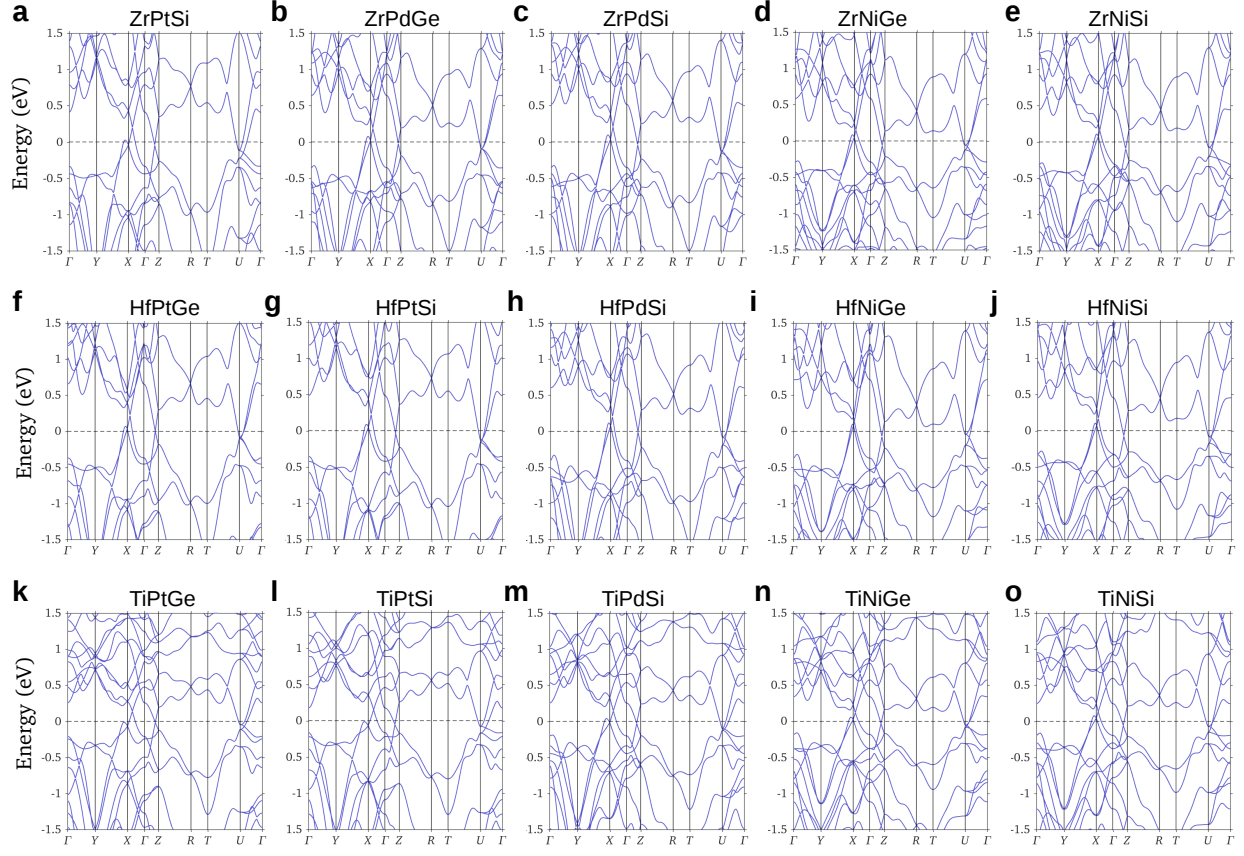


FIG. 11. Bulk band structures of TT'X family (without SOC). (a) ZrPtSi, (b) ZrPdGe, (c) ZrPdSi, (d) ZrNiGe, (e) ZrNiSi, (f) HfPtGe, (g) HfPtSi, (h) HfPdSi, (i) HfNiGe, (j) HfNiSi, (k) TiPtGe, (l) TiPtSi, (m) TiPdSi, (n) TiNiGe, and (o) TiNiSi.

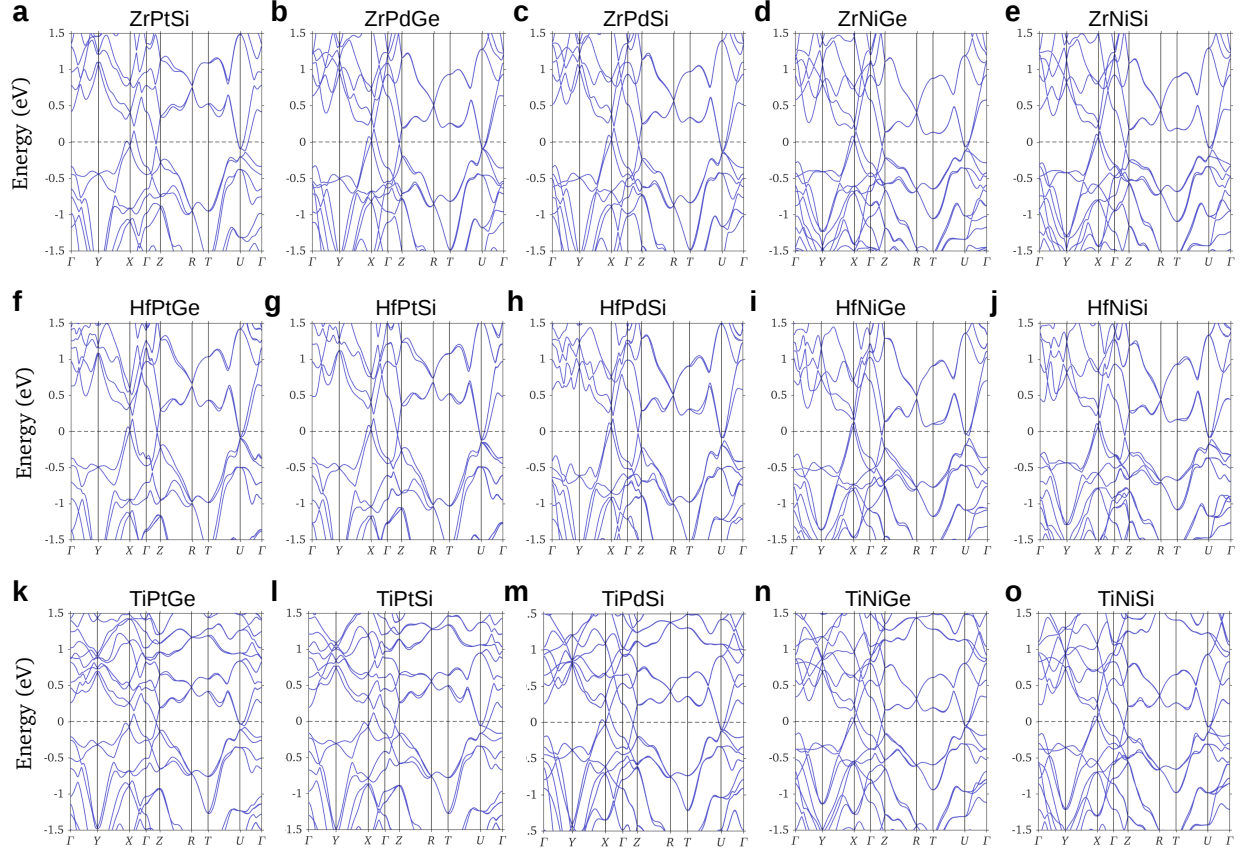


FIG. 12. Same as the caption to Fig. 11, except that this figure refers to band structure with SOC.