A new self-filling mechanism of band gap in magnetically doped topological surface states: spin-flipping inelastic scattering

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Abstract
We investigate the influence of in-plane spin-exchange interactions on a topological insulator (TI) surface doped with nanomagnets under the second perturbation theory. We propose a novel self-filling mechanism of the surface-state band gap. It is found that when the out-of-plane exchange coupling favors an energy gap around the Dirac point, the in-plane component tends to suppress the induced gap, and even fill it completely. Our theory is based on the spin-flipping inelastic scattering, which creates a complex structure of self-energy, effectively modifying the band gap by renormalizing the magnetic moment and chemical potential. We explicitly analyze the filling effect in the electronic dispersion relation and density of states for different scenarios set by systemic parameters. This self-filling effect induced by spin-exchange coupling itself opens new perspectives for understanding of various magnetically doping phenomena on the TI materials and is expected to mediate the controversy concerning the magnetically doping induced gap.

Introduction
Topological insulators (TIs) as a new class of topologically non-trivial states of matter are currently attracting great interest in condensed matter physics due to their promising applications for spintronics and topological quantum computation [1]. One of their remarkable hallmarks is the unique gapless topological surface states (TSSs), which possess the spin-momentum locking nature protecting the Dirac electrons immune from the backscattering off non-magnetic impurities [2–4].

Requirement for the manipulation of the Dirac electronic properties, which is of critical importance for spin-based spintronics, challenges the ability to open up a band gap at the topologically protected Dirac point. Among various techniques, the most effective one is magnetic doping in the bulk [5–8] or on the surface [9–11] of TIs since the magnetic impurities are expected to generate an effect magnetic field. There are tremendous efforts carried out along this line. However, up to now there still is no consensus on the behavior of magnetic impurities experimentally and theoretically. When some works [6–9, 12, 13] reported that magnetic doping can destroy the Dirac point of the TSSs by opening a local gap, some others [10, 11, 14–18] demonstrated that the Dirac node remains immune from magnetic perturbations. For instance, in Fe-deposited Bi$_2$Se$_3$ materials, an energy gap with almost 100 meV was reported in the bulk [8] or surface [9] doping, but later researches [17] found no energy gap and instead a strong resonance around the Dirac point, revealing a surprising robustness of the TSS against magnetic moments. Subsequent several experiments [11, 18] also supported the conclusion of no gap opening. Especially, some studies [10, 19] even showed that there is essentially no difference between magnetic and non-magnetic impurities.

How to understand these contradictory results becomes interesting and challenging. Many studies concentrate on the formation of long-range ferromagnetic order in Ruderman–Kittel–Kasuya–Yosida (RKKY) mediation [9, 13–15, 20] and the screening effect of impurity spin induced by the Kondo effect [21, 22]. Recently, Black-Schaffer et al. [23] suggested that the electron scattering off a strong non-magnetic potential is a promising mechanism to fill the energy gap, trying to reconcile these conflicting claims. Meanwhile, they also
theoretically predicted the opposite scenario that the non-magnetic impurities with spatially extended structure can cause gap-like opening [24] by skew scattering. In realistic experiments, we notice that the most magnetic impurities absorbed on TI surfaces such as Fe and Co possess a large in-plane magnetic anisotropy [11, 17, 18], leading to strong in-plane spin-exchange coupling between the impurities and the topological surfaces. In previous theoretical studies, however, one can notice that the contribution from the in-plane component is usually ignored since it in mean-field theory only shifts the position of dispersion relation in momentum space and is gauged out [25].

In this paper, we propose a novel self-filling mechanism of the band gap. By emphasizing the role of in-plane spin-exchange coupling of surface electrons to magnetic impurities, we study the spin-inelastic scattering within the second-order perturbation theory. In fact, the spin-inelastic scattering induced by the in-plane spin-exchange interactions is extensively investigated for the nanomagnets absorbed on the conventional substrate or single molecular junction [26–29], but it is received much less attention in TIs [30, 31]. We here find that the spin-inelastic scattering can create the complex energy dependence of self-energy. When the out-of-plane exchange opens a band gap around the Dirac point, the in-plane component effectively suppressing the band gap and even completely fill it up.

Model and method

We consider a TI surface doped with high-spin quantum magnets, e.g., magnetic atoms or magnetic ions. The magnetic anisotropy of the quantum magnet \( S = (S_x, S_y, S_z) \) is described with spin Hamiltonian

\[
H_{\text{spin}} = -DS_z^2 + E(S_x^2 - S_y^2),
\]

where \( D \) and \( E \) define the uniaxial and transverse anisotropy constants, respectively. We assume the quantum magnets to be point-like spin \( S(\mathbf{r}) = \sum_{\mathbf{r}_i} S(\mathbf{r}_i) \delta (\mathbf{r} - \mathbf{r}_i) \) where \( \delta (\mathbf{r} - \mathbf{r}_i) \) is a delta function, randomly distributed at the spatial position \( \mathbf{r}_i \). They interact with the surface electrons in the form of the Heisenberg spin exchange [13]

\[
H_{\text{ex}} = H_{\text{ex}}^{\parallel} + H_{\text{ex}}^{\perp},
\]

Here, \( S(\mathbf{r}) = \frac{1}{2} \epsilon \ell_i(\mathbf{r}) \sigma(\mathbf{r}) \) is the spin of conducting electrons, expressed with the vector of Pauli matrices \( \sigma \) and the creation operator of electrons \( \epsilon(\mathbf{r}) = \sum_{\mathbf{k}} e^{-i \mathbf{k} \cdot \mathbf{r}} \hat{c}_{\mathbf{k} \sigma}^\dagger \delta(\mathbf{r} - \mathbf{r}_i) \) where spinor \( \hat{c}_{\mathbf{k} \sigma}^\dagger = (\hat{c}_{\mathbf{k} \uparrow}^\dagger, \hat{c}_{\mathbf{k} \downarrow}^\dagger) \) and planar wave vector \( \mathbf{k} = (k_x, k_y), J_\parallel (J_\perp) \) is the anisotropic spin-exchange constant perpendicular (parallel) to the TI surface. The \( J_\parallel \) term, describing the spin-conserving elastic scattering, is responsible for the ferromagnetic ordering. At the mean-field level, which is valid for not too strong exchange, it can be written as \( H_{\text{ex}}^{\parallel} = M_z \sigma_z \) in terms of a collective moment \( M_z = \frac{n_i J_{\text{ex}}}{k_B T_m} (S_z) \) where \( n_i \) is average impurity concentration. Here, a finite \( M_z \) is assumed as a result of spontaneous ferromagnetic order of impurity spins (e.g., RKKY interaction) for sufficient surface coverage or an applied field. Under the mean-field approximation, \( H_{\text{ex}}^{\parallel} \) can be separated from \( H_{\text{ex}}^{\perp} \) and added to the Hamiltonian of TI surface, given by

\[
H_{\text{TI}} = \hbar v_F(k_x \sigma_x + k_y \sigma_y) + M_z \sigma_z - \mu.
\]

Obviously, the dispersion of \( H_{\text{TI}} \), \( \epsilon_{\mathbf{k}} = \sqrt{(\hbar v_F k)^2 + M_z^2} - \mu \) with the chemical potential \( \mu \), opens an energy gap of \( 2M_z \) at the Dirac point between \( \gamma = + \) conduction and \( \gamma = - \) valence bands.

On the other hand, the in-plane spin-exchange coupling \( H_{\text{ex}}^{\perp} \) is responsible for the spin-flipping inelastic scattering. To include its contribution, we calculate the correction to the electron self-energy within second-order perturbation theory following the standard procedure [32] and assuming the localized spin changing slowly compared with the electron transport relaxation time, the impurity averaged self-energy matrix to first order in the density of impurities is derived as [30, 31]

\[
\Sigma(\omega_n) = -n_i \left( \frac{k_B T_m}{2\hbar v_F} \right)^2 \sum_{\mathbf{k}, \mathbf{k}'} \sum_{M, M'} (\sigma \cdot S_{MM'}) G^0(\mathbf{k}, \mathbf{k}'\omega_n - i\nu_0)
\]

\[
\times (\sigma \cdot S_{MM'}) D_M^{\sigma}(\omega_n + i\nu_0) D_{M'}^0(\omega_n).
\]

Here, we have used the Green’s function with respect to equation (1), \( G^0(\mathbf{k}, \mathbf{k}\omega_n) = 1/[(\mathbf{k}\omega_n - H_{\text{TI}})] \), the Fourier transform of the pseudofermion Green’s function [30] \( D_M^{\sigma}(\tau, \tau') = -i \langle T M_{\sigma}(\tau) d_M(\tau') \rangle \), defined with quasiparticle operators in the eigenstates \( |M\rangle \) of equation (1), and \( S_{MM'} = \sum_{M''} S_{MM''} d_M^{\dagger}(\tau) d_{M''}(\tau) \) with \( S_{MM'} = (M'[S]|M) \). Using Matsubara frequencies \( \omega_n = (2n + 1)\pi k_B T \) for fermions and \( \nu_0 = 2n\pi k_B T \) for bosons, we calculate the self-energy in equation (4) as
with a scalar part $\Sigma_0(i\omega_n) = \sum_{i} B_{k_i}(i\omega_n)$ and magnetic part $\Sigma_2(i\omega_n) = M_z \sum_{i} B_{k_i}(i\omega_n)/\xi_{k_i}$, where

$$B_{k_i}(i\omega_n) = \frac{n_{f,0}^2}{8(\hbar v_F)^2} \sum_{M,M'} (|S_{M,M'}^+|^2 + |S_{M,M'}^-|^2)(P_M - P_{M'}) \frac{1 + n_s(E_{M'M}) - f(i\omega_n)}{i\omega_n - \mu - \xi_{k_i} - E_{M'M}}.$$ (6)

Here, $E_{M'M} = E_{M'} - E_{M}, P_M = \exp(-E_M/k_B T)/\sum_{i} \exp(-E_i/k_B T), E_M$ is the eigenvalue of $H_{\text{spin}},$ and $f(x)$ ($n_{f,0}(x)$) is the Fermi (Bose) distribution functions. The dynamic self-energies $\Sigma_0(i\omega_n)$ and $\Sigma_2(i\omega_n)$ are complex and energy-dependent, carrying the information about magnetic structure of nanomagnets through the common factor $B_{k_i}(i\omega_n)$, which describes the localized spin moment undergoing spin-flipping transitions ($S_{M,M'}^+$ or $S_{M,M'}^-$) between magnetic states with different energies ($E_{M'} = E_{M'}$). These processes directly change the $\xi_0$ part and $\xi_{k_i}$ part in $G^0(k, i\omega_n)$.

With the spin-flipping contribution, the full fermion Green’s function is corrected by Dyson equation

$$G(k, i\omega_n) = G^0(k, i\omega_n)[1 - \Sigma(i\omega_n)]G^0(k, i\omega_n)]^{-1}$$ and given by

$$G(k, i\omega_n) = \frac{1}{2} \sum_{\gamma=\pm} [1 + \gamma \mathbf{n}_k \cdot \sigma] G_{\gamma}^0(k, i\omega_n),$$ (7)

where we introduce the band-dependent Green’s function $G_{\gamma}^0(k, i\omega_n) = [i\omega_n - \Sigma_0(i\omega_n) + \mu - \gamma \xi_{k}]^{-1}$ and effective momentum unit vector $\mathbf{n}_k = (\hbar v_F k_x, \hbar v_F k_y, M_z + \Sigma_2(i\omega_n))/\xi_{k}$ with $\xi_{k} = \sqrt{(\hbar v_F k)^2 + [M_z + \Sigma_2(i\omega_n)]^2}$. Note that $\Sigma_2(i\omega_n)$ directly renormalizes the gap $M_z$ while $\Sigma_0(i\omega_n)$ modifies the chemical potential $\mu$, both of which will shrink the existing energy gap but in different means as discussed below.

**Results and discussion**

**Filling of gap for $\mu < M_z$**

The non-trivial dependence of self-energy on chemical potential $\mu$ will lead to $\mu$-resolved electron spectrum structure. We first discuss the case of $\mu$ within the band gap. In figures 1(a)–(d) we display the evolution of the band gap in spectral function $A(k, \omega) = -\frac{1}{\pi} \text{Im} \text{Tr} G(k, i\omega_n - \omega + i\delta)$ for increasing in-plane coupling $J_{||}$ with magnetic impurities. Without the coupling $J_{||} = 0$, a distinct band gap appears, centered around the Dirac point, due to the massive Dirac dispersion $\varepsilon_{k_x} = \gamma [(\hbar v_F k)^2 + M_z^2]^{1/2}$ while a linear dispersion retains far away from the Dirac point. As $J_{||}$ is enhanced gradually, two edges of the band gap move towards the Dirac point,
making the gap narrower and narrower and finally disappeared completely, see figure 1(d). This indicates that the self-energy due to correction of spin-inelastic scattering tends to fill the energy gap.

Above behaviors can be understood well from the structure of dynamic self-energy. At zero temperature, we can approximate \( f(x) = \Theta(\mu - x) \) as a Heavisider step function, \( \eta_0(x > 0) = 0 \), and occupations \( P_{\pm S} = 1 \) and zero otherwise. Thus, in equation (6) only need to consider the inelastic transitions involving the first excited state and the ground state, i.e., \( \pm(S - 1)|S_\pm|\pm S \sqrt{2S} \), and neglect other transitions between high excited states safely. Finally, we obtain a simple analytical form of the self-energy as

\[
\Sigma_2(\omega) = \eta M_z \left[ \ln \frac{M_z - \omega^-}{E_c - \omega^-} + \ln \frac{M_z + \omega^+}{E_c + \omega^+} \right] - i\pi \Theta(\omega^- - M_z) + i\pi \Theta(-\omega^+ - M_z),
\]

where \( \eta = \frac{n_{gS}}{\pi(\hbar/2)^2} \), \( E_c \) is the cutoff energy for the band width of surface states, and \( \omega^\pm = \tilde{\omega} \pm \Delta \) with \( \tilde{\omega} = \omega - \mu \) and the energy threshold \( \Delta = |E_{\pm S, \pm S} - \mu| \) between ground states \( |\pm S \rangle \) and first excited states \( |\pm S \pm 1 \rangle \). \( \Sigma_0(\omega) \) has the same expression but adding an extra energy factor \( \omega^\pm \) before each term to change the weight. We plot the resulting \( \Sigma_2(\omega) \) and \( \Sigma_0(\omega) \) as a function of energy in figures 1(e) and 1(f), respectively. Within the band gap \( [-M_z, M_z] \), both the imaginary parts vanish due to energetically prohibited magnetic excitations and thus the contribution to gap closing is attributed completely to the real parts of dynamic self-energies. Importantly, \( \text{Re}[\Sigma_2(\omega)] \) is always negative in this interval, so tending to reduce magnetic gap owing to its renormalized effect \( M'_z + \Sigma_2(\omega) \) in equation (7). By contrast, \( \text{Re}[\Sigma_0(\omega)] \) modifies the band gap in a different way. In figure 1(f), \( \text{Re}[\Sigma_0(\omega)] \) is antisymmetric with respect to \( \omega \), negative for \( \tilde{\omega} > 0 \) while positive for \( \tilde{\omega} < 0 \). This interesting energy-dependent self-energy pushes the valence band edge to high energy and the conducting band edge to low energy, thus narrowing the band gap. The present physics of gap closing below the magnetic excitations and the broadening of spectral function in the region of \( |\omega| > M_z + \Delta \) is attributed to the finite imaginary part of self-energies, which can be clearly visible from the self-energy in figures 1(e) and 1(f).

In figures 2(a) and (b), we display the variation of density of states (DOSs), defined as \( \rho(\omega) = -\frac{1}{\pi} \text{Im} \Sigma_2 \text{TrG}(\mathbf{k}, i\omega_n \rightarrow \omega + i\delta) \), at zero and finite temperature, respectively. At zero temperature as shown in

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure2.png}
\caption{Variation of DOS for different \( J_z \) at zero (a) and finite temperature \( k_B T = 1 \) (b). (c) Dependence of renormalized moment \( M'_z \) on \( J_z/J_\perp \) and (d) DOS \( \rho(\omega) = 0 \) at the Dirac point changes with finite temperature. The other parameters are the same as in figure 1.}
\end{figure}
figure 2(a), with increasing \( J_0 \) the effect of gap filling reappears. The minimum \( J_0 \) value to completely fill the gap is determined by \( M_z^{\text{eff}} = 0 \) with \( M_z^{\text{eff}} = M_z + \text{Re} \{ \Sigma_z(\omega) \} \), corresponding to the intersection point of curves in figure 2(c), where \( M_z^{\text{eff}} \) is plotted as a function of \( J_0/J_0^* \). Note that \( \text{Re} \{ \Sigma_z(\omega) \} \) is not important for determining the critical value of \( J_0 \) since it is almost zero around \( \omega = 0 \), despite suppressing the gap far away from the Dirac node. Two sharp peaks of magnetic excitations, pinned at \( \omega = \pm (M_z + \Delta) \), is remarkable, originating from the logarithmic term in equation (8). At finite temperature, we perform the numerical calculations and the results are presented in figure 2(b), with the same other parameters. By contrast with figure 2(a), besides the magnetic excitations are smoothed out, the most prominent distinction is that the perfect filling effect becomes easier, even before \( J_0 \) reaching the zero- \( T \) threshold. With increase in \( J_0 \), the gapped Dirac cone forms a flatten dip structure, instead of V-shape, and the DOS at Dirac point is overall lifted upwards. We in figure 2(d) plot \( \rho(\omega = 0) \) as a function of temperature for different \( J_0 \), where the DOS at Dirac point increases with \( T \) and quickly reaches a saturate plateau with height proportional to \( J_0 \). All the distinctions at finite temperature are due to the appearance of the imaginary part of self-energies \( \Sigma_z(\omega) \) and \( \Sigma_0(\omega) \), which not only renormalize the band gap as \( M_z^{\text{eff}} = \sqrt{(M_z + \text{Re}[\Sigma_z(\omega)])^2 + (\text{Im}[\Sigma_z(\omega)])^2} \), but also cause finite lifetime of quasi-particles. Physically, the temperature facilitates the inelastic scattering transitions between spin states by broadening the magnetic levels of impurities. Finally, we want to remark that the gap filling and magnetic excitations are unique for nanomagnets due to their magnetic anisotropy. If the magnetic anisotropy constants \( D \) and \( E \) vanish, above features are lost due to \( \Sigma(i\omega_0) = 0 \), which can be seen from equation (6) where \( P_M = P_{M'} \). Finite \( D \) or \( E \) is necessary to create the level difference. Both of them affect the spectral function and DOS through entering eigen–energy \( E_M \). With increasing of \( D \) or \( E \), the magnetic excitation peaks, located at \( \omega = \pm (M_z + \Delta) \), move away from the Dirac point, and meanwhile the magnitude of self-energy decreases because large level difference \( E_{MM'} \) reduces the spin–inelastic scattering processes. Thus, we need larger \( J_0 \) determined by \( M_z^{\text{eff}} = 0 \), to fill the gap. In realistic TI systems doped magnetically, we take the typical magnetic parameters \( D = 1.33 \text{ K}, E = -0.034 \text{ K}, \) and \( S = 4 \) for nanomagnets[33], and the typical parameters \( Ec = 300 \text{ meV}, h\nu_F = 3.7 \text{ eV Å}, k_BT = 0.3 \text{ K}, n_i = 0.1, \) and \( M_z = 20 \text{ meV} \) for TIs[17]. We estimate that the minimum value to completely close the energy gap is \( J_0 = 0.74 \text{eV} \). If the temperature rises to be \( k_BT = 10 \text{ K} \), the critical value reduces to be \( J_0 = 0.51 \text{eV} \). The in-plane overlapping values between the TI surface and the magnetic impurities are within the experimentally relevant range.

Filling of gap for \( \mu > M_z \)
The above discussions are for the chemical potential set within the band gap. If \( \mu > M_z \) out of the gap, the scenario is very different. For convenience, we start with the self-energy structure. At zero temperature, we can derive its analytical form as

\[
\Sigma_z(\omega) = \eta M_z \left\{ \ln \left| \frac{M_z - \mu}{\omega^+ - \mu} \right| - \ln \left| \frac{E_c - \mu}{\omega^- - \mu} \right| - \ln \left| \frac{E_c + \omega^+}{M_z + \omega^+} \right| - \text{i} \pi \Theta (-\omega^- + \mu) \Theta (\omega^+ - M_z) \right. \\
- \left. \Theta (-\omega^- - M_z) - \text{i} \pi \Theta (\omega^- - \mu) \right\}. \quad (9)
\]

\( \Sigma_z(\omega) \) and \( \Sigma_0(\omega) \) are depicted in figures 3(c) and (f). Different from equation (8), the real parts of both \( \Sigma_z(\omega) \) and \( \Sigma_0(\omega) \) have four magnetically excited peaks, respectively, located at \( \omega = \mu \pm \Delta \) and \( \omega = -\Delta \pm M_z \), exhibiting complex dependence on energy. Similar to the discussions for \( \mu < M_z \), negative \( \text{Re} \{ \Sigma_z(\omega) \} \) and asymmetric \( \text{Re} \{ \Sigma_0(\omega) \} \) guarantee the gap reduction effectively. This point can be seen by comparing the energy gap of the spectral function in figures 3(a) with (b)-(d), where the gap is only partly filled because of setting large moment \( M_z = 0.1 \) and small coupling \( J_0 = 0.45 \). Most interestingly, it is noted that the imaginary parts of both self-energies are non-zero in a central region \( [M_z - \Delta, \mu - \Delta] \) even at zero temperature. This special energy dependence of self-energy generates a profound effect on the corresponding spectral function \( A(k, \omega) \), namely, a horizontal fuzzy strip across the dispersion as shown in figures 3(b)-(d). The central position of the horizontal strip can be tuned by magnetic anisotropy parameter \( D \) or \( E \) of nanomagnet while its width is determined by \( \mu \). Therefore, tuning \( D \) or \( E \) can make the horizontal strip move from conducting band to valence band with unchanged shape, resembling to the behavior of one type of quasi-particle excitations. Importantly, when the horizontal band is shifted into the band gap as depicted in figure 3(c), the gap is filled partly. If one increases the chemical potential \( \mu \) further, the horizontal band becomes broader and broader and finally fill the gap completely at the Dirac point, as illustrated in figure 3(d). Physically, it is a consequence of strongly mixing between Dirac electrons and magnetic excitations.
In figures 4 (a) and (b) we present the variation of DOS with energy for chemical potential $\mu$ out of the band gap, with the parameters corresponding to figures 3 (c) and (d), respectively. Again, a remarkable bump with width of $|\mu-M_z|$ emerges within the band gap. For small $\mu$, the bump partly fills the band gap and at its edge forms two DOS dips, resembling to the splitting of Dirac point [34]. For large $\mu$ as in figure 4 (b), the broaden bump almost fill the gap. The numerical results at finite temperature are also shown in figure 4, indicated with line-symbols, where the gap of the DOS at the Dirac point is suppressed perfectly by finite temperature.

**Summary**

We have investigated the spin-inelastic scattering of Dirac electrons off anisotropic nanomagnets doped on a TI surface. We mainly focus on the role of the in-plane spin exchange coupling of Dirac electrons to nanomagnets, responsible for spin flipping, and find it can significantly modify the topologically protected Dirac electron spectrum by creating complex energy-dependence of self-energy. For gapped Dirac dispersion induced by out-
of-plane spin exchange, the dynamic self-energy can modify the band gap by directly renormalizing the magnetic moment and the chemical potential, and even by introducing new spectrum structures. Consequently, the band gap around the Dirac point in the dispersion spectrum as well as the DOS can be filled completely or partly, depending on specific parameters, for which we have given an explicit analysis. This self-filling mechanism induced by the spin-exchange coupling itself opens new perspectives for understanding the magnetically doping effect on TI materials.

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References

[1] Hasan M Z and Kane C L 2010 Rev. Mod. Phys. 82 3045
[2] Qi X L and Zhang S C 2011 Rev. Mod. Phys. 83 1057
[3] Roushan P, Seo J, Parker C V, Hor Y S, Hsieh D, Qian D, Richardella A, Hasan M Z, Cava R J and Yazdani A 2009 Nature 460 1106
[4] Beidenkopf H, Roushan P, Seo J, Gorman L, Drouid I, Hor Y S, Cava R J and Yazdani A 2011 Nat. Phys. 7 939
[5] Vobornik I, Manju U, Fujii J, Borgatti F, Torelli P, Krizmanic D, Hor Y S, Cava R J and Panaccione G 2011 Nano Lett. 11 4079
[6] Checkelsky J G, Ye J, Onose Y, Iwasa Y and Tokura Y 2012 Nat. Phys. 8 729
[7] Xu S Y et al 2012 Nat. Phys. 8 616
[8] Chen Y L et al 2010 Science 329 659
[9] Wray L A, Xu S Y, Xia Y, Hsieh D, Fedorov A V, Hor Y S, Cava R J, Bansil A, Lin H and Hasan M Z 2011 Nat. Phys. 7 32
[10] Valla T, Pan Z H, Gardner D, Lee Y S and Chu S 2012 Phys. Rev. Lett. 108 117601
[11] Scholz M R, Sánchez-Barriga J, Marchenko D, Varykhalov A, Volykhov A, Yashina I V and Rader O 2012 Phys. Rev. Lett. 108 256810
[12] Abanin D A and Pesin D A 2011 Phys. Rev. Lett. 106 136802
[13] Liu Q, Liu C X, Xu C, Qi X L and Zhang S C 2009 Phys. Rev. Lett. 102 156603
[14] Biswas R R and Balatsky A V 2010 Phys. Rev. B 81 233405
[15] Sessi P, Reis F, Bathon T, Kokh K A, Tereshchenko, O E and Bode M 2014 Nat. Commun. 5 5349
[16] Henk J, Ernst A, Eremeev S V, Chulkov E V, Maznichenko I V and Mertig I 2012 Phys. Rev. Lett. 108 206801
[17] Honolka J et al 2012 Phys. Rev. Lett. 108 256611
[18] Schlenk T et al 2013 Phys. Rev. Lett. 110 126804
[19] Bianchi M, Hatch R C, Mi J, Iversen B B and Hofmann P 2011 Phys. Rev. Lett. 107 086802
[20] Schmidt T M, Miwa R H and Fazzio A 2013 J. Phys.: Condens. Matter 25 445003
[21] Orignac E and Burdin S 2013 Phys. Rev. B 88 035411
[22] Yu Y H, She L M, Fu H X, Huang M, Li H, Meng S and Cao G Y 2014 ACS Nano 8 11576
[23] Black-Schaffer A M, Balatsky A V and Fransson J 2013 Phys. Rev. B 88 035411
[24] Fransson J, Black-Schaffer A M and Balatsky A V 2014 Phys. Rev. B 90 241409(R)
[25] Parhiizar F, Moghaddam A G and Asgari R 2015 Phys. Rev. B 92 045429
[26] Heersche H B, de Groot Z, Folk J A, van der Zant H S J, Romeike C, Wegewijs M R, Zobbi L, Barreca D, Tondello E and Cornia A 2006 Phys. Rev. Lett. 96 206801
[27] Kim G H and Kim T S 2004 Phys. Rev. Lett. 92 137203
[28] Oberg J C, Calvo M R, Delgado F, Moro-Lagares M, Serrate D, Jacob D, Fernandez-Rossier J and Hirjibehedin C F 2013 Nat. Nanotechnol. 8 4
[29] Wang R Q, Sheng L, Shen R, Wang B G and Xing D Y 2010 Phys. Rev. Lett. 105 057202
[30] Thalmeier P and Akbari A 2012 Phys. Rev. B 86 245426
[31] Wang R Q, Sheng L, Yang M, Wang B G and Xing D Y 2015 Phys. Rev. B 91 245409
[32] Mahan G D 1990 Many-Particle Physics 2nd edn (New York: Plenum)
[33] Leuenberger M N and Muccio E R 2006 Phys. Rev. Lett. 97 126601
[34] Black-Schaffer A M and Balatsky A V 2012 Phys. Rev. B 85 121103(R)