Surface spectral function of momentum-dependent pairing potentials in a topological insulator: application to Cu$_x$Bi$_2$Se$_3$

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Abstract

We propose three possible momentum-dependent pairing potentials as candidates for topological superconductors (for example Cu$_x$Bi$_2$Se$_3$), and calculate the surface spectral function and surface density of states with these pairing potentials. We find that the first two can give the same spectral functions as the fully gapped and node-contacted pairing potentials given by Fu and Berg (2010 Phys. Rev. Lett. 105 097001), and that the third one can obtain a topological non-trivial case in which there exists a flat Andreev bound state and which preserves the threefold rotation symmetry. We hope our proposals and results will be assessed by future experiment.

(Some figures may appear in colour only in the online journal)

Recently, topological insulators [1–3] have attracted great attention in condensed matter physics for their physical properties. The experimental [4–9] and theoretical [10–15] investigations have shown that topological insulators (TIs) are fully gapped in the bulk but gapless on the surface, which is protected by time-reversal symmetry. These surface states indicate massless Dirac fermions. The research on topological insulators was then soon generalized to topological superconductors [16–18]. Similarly, research shows that topological superconductors also have gapless surface states, which indicate massless Majorana fermions [19], and that the property of three-dimensional topological superconductivity is protected by particle–hole and time-reversal symmetry and characterized by a topological invariant [16, 20]. As we know, the Majorana fermions are of great interest in fundamental physics and have potential applications in quantum computation [21, 22].

Experiment [23] shows that a superconductive phase is induced at transition temperature $T_c = 3.8$ K when copper atoms are doped into the topological insulator Bi$_2$Se$_3$ with the concentration of copper in the range 0.12–0.15, and the work [24] shows that the surface state of Cu$_x$Bi$_2$Se$_3$ is topologically non-trivial. Recent experiments [25, 26] further confirm the existence of topological surface states by measuring the zero-bias conductance peak (ZBCP). In the low-transmissivity limit [27], the ZBCP corresponds to the surface density of states (SDOS) and indicates an unconventional Andreev bound state (ABS) [28, 29] on the surface. However, a more detailed analysis in the supplemental material of [25] shows that a gapless and topologically non-trivial bulk band structure may be preferred.

Up to now, the exact pairing mechanism is still unclear, but there are some theoretical proposals on the pairing symmetry of Cu$_x$Bi$_2$Se$_3$, including the odd-parity pairing potential proposed by Fu and Berg [30] and Sato [31, 32], and the explanation of Sasaki et al [25] according to experimental results. We analyzed all four kinds of momentum-independent pairing potential and found that it is impossible to construct a proper pairing potential which satisfies the following conditions: (1) it is topologically non-trivial, (2) the bulk gap
has point nodes, and (3) the pairing potential is invariant under symmetric operations of the D$_{3d}$ group. The first two constraints come from experimental results and the last constraint is caused by the symmetry properties of the Bi$_2$Se$_3$ crystal. Theoretical investigations on momentum-dependent pairing potentials are inevitable.

In this paper, we propose three momentum-dependent pairing potentials for Cu$_2$Bi$_2$Se$_3$, and calculate the surface spectral functions with these pairing potentials. We find that of our pairing potentials the first two sorts are similar to $\Delta_2$ and $\Delta_4$ proposed in [25] and can obtain what $\Delta_2$ and $\Delta_4$ give, and that the third one can give a topologically non-trivial case in which there exists a flat ABS and which preserves the threefold rotation symmetry.

As reported [24], near the $\Gamma$-point, the band dispersion of the normal state of Cu$_2$Bi$_2$Se$_3$ can be described by the low energy effective Hamiltonian [14, 15] for Bi$_2$Se$_3$, with a finite chemical potential in the conduction band induced by copper doping. The Hamiltonian is

$$ h(k) = (M \tau_z - \mu) + \tau_x \left( A(k_x \sigma_x + k_y \sigma_y) + B k_c \sigma_z \right), $$

where $M$ is the rest mass, $\mu$ is the chemical potential, $A$ and $B$ are Fermi velocities along different directions, $\tau_z = \pm 1$ denote the two orbits, and $\sigma_{x,y,z}$ are spin Pauli matrices. For the superconductivity phase, the Hamiltonian can be written in the Bogoliubov–de Gennes (BdG) formalism:

$$ \mathcal{H} = \begin{pmatrix} h(k) & \Delta(k) \\ \Delta^\dagger(k) & -h^T(-k) \end{pmatrix}, $$

where $\Delta(k)$ is the pairing potential ($4 \times 4$ matrix). For the time-reversal invariant case, the pairing potential can be classified into two categories according to inversion symmetry:

$$ \Delta_1(k) = i k_y \sigma_y \left( \sigma_y \Delta_{1,aa}^{aj} + \tau_z \sigma_y \Delta_{1,as}^{aj} \right), $$

$$ \Delta_2(k) = i k_y \sigma_y \left( \tau_y \sigma_y \Delta_{2,aa}^{aj} + i \tau_z \Delta_{2,as}^{aj} \right), $$

where $\Delta_{1,aa}^{aj}, \Delta_{1,as}^{aj}, \Delta_{2,aa}^{aj}$ and $\Delta_{2,as}^{aj}$ are real functions of momentum and inversion symmetric; the summation convention is used in this paper. One can find that $\Delta_2(k) = \tilde{P}^{-1} \Delta_2(k) \tilde{P}$ is inversion symmetric and $\Delta_1(k) = -\tilde{P}^{-1} \Delta_1(k) \tilde{P}$ is inversion anti-symmetric, where the inversion operator is $\tilde{P} = \tau_z$ in coordinate space. As we know, if the pairing potential is dominated by $\Delta_1(k)$ and fully gapped in the Brillouin zone, the criterion for a topological odd-parity superconductor [30–32] claims that the system is topologically non-trivial.

In the following, we consider three typical cases of $\Delta_1(k)$ and calculate the surface spectral functions with them. In order to calculate the spectral functions numerically, we consider a simplified lattice model [25, 33, 34], of which the low energy effective Hamiltonian is equation (1). For the normal state of the Hamiltonian, we use the model and parameters given in the supplemental material of [25].

In the first case, for inversion-anti-symmetric superconductive potential (3), we consider $\Delta_{1,as}^{aj} = 0$ and

$$ \Delta_{1,s} = \begin{pmatrix} A & 0 \\ 0 & B \end{pmatrix}, $$

as an example, and other cases are similar. In this case, the pairing potential takes the form

$$ \Delta(k) = \Delta \begin{pmatrix} A(k_y + i k_z) & B k_z \\ B k_z & -A(k_y - i k_z) \end{pmatrix} \otimes \tau_0, $$

where $\Delta$ is a dimensionless parameter determined by the energy gap of superconductivity, and $\tau_0$ is a $2 \times 2$ identity matrix in orbital space. One of the compactifications of equation (6) can be given as

$$ \Delta(k) = \Delta \begin{pmatrix} A_{1}^{\tau} \ A_{1}^{\tau} \\ A_{1}^{\tau} \ A_{1}^{\tau} \end{pmatrix} \otimes \tau_0, $$

where we refer to the definition of $A_{1}^{\tau}, A_{2}^{\tau}$ in the supplemental material of [25], and take the same parameter values as [25] in our numerical calculation. In this case, the lattice model can preserve the same translation symmetry of the discrete version of $h(k)$ and turn to equation (6) in the low energy limit ($k \rightarrow 0$). We must point out that this lattice model to reproduce the low energy effective Hamiltonian is not valid for $k \gg k_F$.

By using the method in [33, 34], we can obtain the surface spectral function. Considering a semi-infinite system which has a surface at $z = 0$, the momentum which is parallel to the surface $k_\parallel = (k_x, k_y)$ is a good quantum number, and the partition function of the system with an open surface at $z = 0$ can be written as

$$ Z = \int D\psi D\psi^\dagger \exp \left\{ i \int d\tau \sum_{k_\parallel} \sum_{n=0}^{\infty} \left[ \psi_n^\dagger (i \partial_\tau - H_{k_\parallel}) \psi_n + \left( \psi_n^\dagger V_{k_\parallel} \psi_{n+1} + \text{h.c.} \right) \right] \right\}, $$

where $\psi_n$ is the wavefunction for the $n$th layer, $H_{k_\parallel}$ is the intralayer Hamiltonian, $V_{k_\parallel}$ is the interlayer coupling, and h.c. means Hermitian conjugate. The recursive integration layer by layer gives the following retarded Green’s function for the surface state:

$$ G^{-1}(k_\parallel, \omega) = G_0^{-1}(k_\parallel, \omega + i 0^+) - V_{k_\parallel} G^{-1}(k_\parallel, \omega) V_{k_\parallel}, $$

where $G_0(k_\parallel, \omega + i 0^+) = (\omega + i 0^+ - H_{k_\parallel})^{-1}$ is the retarded Green’s function for an isolated layer. The Green’s function of a surface state is calculated by the quick iterative scheme [35] for the transfer matrix. Finally, the surface spectral function is given in the following form:

$$ A(k_\parallel, \omega) = -\frac{1}{\pi} \text{Im Tr } G(k_\parallel, \omega). $$
Figure 1. Model calculation of bulk (a) and surface (b) spectral function $A(k, \omega)$ for BdG Hamiltonian with pairing potential given in equation (5). (c) Surface spectral function as a function of momentum for $\omega = 0$. (d) Bulk (blue dashed line) and surface (red solid line) density of states (DOS). The false color mappings of $A(k, \omega)$ in (a)–(c) are in arbitrary units. Parameters for model calculation are given in the context.

One can also calculate the SDOS by integrating $A(k, \omega)$ over momentum,

$$\rho_s(\omega) = \int \frac{d^2k}{(2\pi)^2} A(k, \omega). \quad (11)$$

In order to agree with the moment-independent pairing, in the calculation, we set the dimensionless parameter $\Delta$ to be $0.15$, the maximum gap size to be about $0.05 \text{ eV}$, and the truncation of momentum $|k_x| = |k_y|$ to be $0.6a^{-1}$ ($a$ is the lattice space in the xy-plane), because, as pointed out above, the lattice model is not valid for $k \gg k_F$.

Now we consider the surface spectral function for some special pairing symmetries. In the first case, the pairing potential is given in equation (6), which looks like a direct sum of two pairing potentials to describe the $^3\text{He}$–B phase. Because, in a topological insulator, strong spin–orbit coupling between different orbits makes the pairing potentials between different energy bands complicated, we must calculate the topological invariant carefully. We can identify the topological invariant by the criterion for a topological odd-parity superconductor [30–32], which shows that the pairing symmetry of equation (6) is topologically non-trivial, and we can also calculate the winding number directly [16],

$$N_w = \frac{1}{24\pi^2} \int d^3k \epsilon^{ijk} \text{Tr}[Q_k^i \partial_j Q_k^j Q_k^k \partial_k Q_k^i \partial_k Q_k^i], \quad (12)$$

where $Q_k = 2P_k - 1$, $P_k = \sum_{n \in \text{occ}} |u_n(k)\rangle \langle u_n(k)|$ is the projector onto the occupied Bloch states. We use the latter method and find that the winding number is totally determined by the topology of the Fermi surface and $N_w = -1\text{Sgn}(\Delta)$ if the pairing potential (6) is non-vanishing only for a thin spherical shell around the Fermi momentum $k \sim k_F$, which implies that, although the pairing potential is written in four bands (different spins and different orbitals), it can be continuously deformed to the weak pairing limit [20] on the Fermi surface. The spectral function for the Hamiltonian with pairing potential (6) (as shown in figure 1) shows that the bulk state is fully gapped (as shown in figure 1(a)) and there is an Andreev bound state on the surface (as shown in figure 1(b)). Similar to the momentum-independent odd-parity pairing potential $\Delta_2$ in [25], this pairing potential is inconsistent with experimental results and there is a minimal of the SDOS at $\omega = 0$ (figure 1(d)), which is not observed at the ZBCP.

One of the explanations indicates that the bulk band structure is topologically non-trivial but with some point nodes, which will induce a flat dispersion of surface helical...
Majorana fermions and contribute a non-vanishing peak of SDOS at zero energy. Among all sorts of momentum-dependent pairing symmetry, there exist some species which possess this property, resembling the momentum-independent odd-parity pairing potential $\Delta_4$ in [25]. We consider the pairing potential given by equation (3) with

$$\Delta_{1,s} = \Delta \begin{pmatrix} 0 & A & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -B \end{pmatrix},$$

and $\Delta_{1,as} = 0$ as another example, and find that the bulk bands have two point nodes at $k = (\pm \sqrt{\mu^2 - M^2}/A, 0, 0)$, and the matrix $Q_k$ for this pairing potential is well defined in the Brillouin zone excluding these two points. In the weak pairing limit, $Q_k$ can be written as

$$Q_k = -i \text{sgn}(\Delta) \left( \begin{array}{ccc} \kappa & \cos \theta \sigma_z & -\sin \theta \sigma_x \\ \kappa & \sin \theta \sigma_x & \cos \theta \sigma_z \\ 0 & 0 & 1 \end{array} \right) \otimes (\tau_0 - \tau_z),$$

near the two singularity points on the Fermi surface, where $\theta = \arctan(B_{k_z}/(A_{k_y}))$ and $\kappa = M/\mu$. Here, we have made a unitary transformation to express $Q_k$ in the eigenvalue representation of $h(k)$ and $\sigma_z$, and $\tau_0, \tau_z$ in the same form as before but with different meanings. Equation (14) shows that there are two ABSs on the boundary of the $xz$-plane, which is similar to the chiral p-wave superconductor but time-reversal symmetry is unbroken here. As shown in the appendix, in this case, the winding number is well defined even though the bulk gap has nodal points. The spectral function of this pairing potential is given in figure 2. The SDOS has a non-vanishing value at $\omega = 0$ (figure 2(d)).

However, this pairing symmetry seems to be also unlikely, because there are only two point nodes of the bulk bands, which breaks the symmetry of the D$_{3d}$ group. As shown in figure 2(c), the spectral function of the surface state for $\omega = 0$ is not invariant under the threefold rotation operation in the $k_xk_y$-plane. As we know, in the $\Delta_4$ suggested for the pairing symmetry of Cu$_x$Bi$_2$Se$_3$ in [25] there also exists such a deficiency.

Now we discuss the other choices in equations (3) and (4). We assume that $\Delta_{1,s}^{aj}$, $\Delta_{1,as}^{aj}$, $\Delta_2^{aj}$ and $\Delta_2^{bj}$ are constants. The pairing symmetry $\Delta_{1,as}^{aj}$ of the anti-symmetric orbits is similar to the pairing symmetry $\Delta_{1,s}^{aj}$ of the symmetric orbits; we can construct parallel theories for pairing potentials and get similar spectral functions and DOSs. For the fully bulk-gapped systems, we find that they can deform to each other continuously. The difference between $\Delta_{1,as}^{aj}$ and $\Delta_{1,s}^{aj}$ cannot be distinguished by the shape of the spectral function. In addition, after calculating the spectral functions for $\Delta_2(k)$ at the linear order of momentum, we find that the system

![Figure 2. Model calculation of bulk (a) and surface (b) spectral function $A(k, \omega)$ for BdG Hamiltonian with pairing potential given in equation (13). (c) Surface spectral function as a function of momentum for $\omega = 0$. (d) Bulk (black dashed line) and surface (red solid line) DOS. The false color mappings of $A(k, \omega)$ in (a)–(c) are in arbitrary units. Parameters for model calculation are given in the context.](image-url)
is bulk gapless when $\Delta_{2}^{0c} \neq 0$ and others are zero, and the Hamiltonian with only $\Delta_{2}^{0j} \neq 0$ can be bulk gapped and topologically trivial; its winding number $N_{W} = 0$. These results are summarized in table 1.

The above analyses show that a pairing potential proportional to the first order of momentum does not provide a proper candidate which satisfies the following conditions: (1) it is topologically non-trivial so that it can lead to a topologically non-trivial ABS, (2) its band structure has some point nodes in the $k_{x}k_{y}$-plane so that it can provide a large SDOS, and (3) the Hamiltonian and the corresponding spectral function preserve the threefold rotation symmetry. Higher orders of momentum must be considered. The simplest pairing potential which preserves the threefold rotation symmetry takes the form of the one-dimensional irreducible representation of the threefold rotation operation,

$$\Delta(k) = \Delta A^{3}(k_{x}^{3} - 3k_{x}^{2}k_{y})\sigma_{z} \otimes \tau_{0};$$

we calculate the spectrum function and DOS for Cu$_{2}$Bi$_{2}$Se$_{3}$ with such a pairing potential, and we find that in the bulk spectral function there exist complex line nodes, and the surface spectral function has similar shape; both of them contribute minimal DOS at $\omega = 0$. Such a result is not consistent with experimental results. In order to obtain a large net surface DOS, we need to open the line nodes to get point nodes and we construct the following pairing potential:

$$\Delta_{1}(k) = \Delta[Bk_{z} + \lambda A^{2}(k_{y}^{2} - 3k_{z}^{2}k_{y})]\sigma_{z} \otimes \tau_{0}, \quad (15)$$

where $\lambda$ is a parameter and we choose $\lambda = 40eV^{-2}$ in practical calculation; for simplicity we also choose $\Delta_{2}^{0j} = \Delta_{2}(k) = 0$. The term proportional to $Bk_{z}$ is applied to open gaps on the line nodes. The Hamiltonian (2) with pairing potential (15) has six point nodes on the equator:

$$M^{2} + A^{2}(k_{x}^{2} + k_{y}^{2}) = \mu^{2} \quad k_{z}^{2}(k_{x}^{2} - 3k_{z}^{2})^{2} = 0; \quad (16)$$

these point nodes are independent of the exact value of parameter $\lambda$. We must emphasize that the $k_{x}^{3}$ terms are not high order corrections, and they are as important as the linear order terms of momentum for the weak pairing limit. The spectral function and SDOS are given in figure 3; now threefold rotation symmetry is preserved (as shown in figure 3(c)); the net effect of flat helical Majorana fermions induced by six point nodes of bulk bands accumulates a sharp surface spectral function peak around the $\Gamma$-point for $\omega = 0$, and this effect is also manifested in the SDOS (as shown in figure 3(d)).

In summary, we investigate possible momentum-dependent pairing potentials in superconducting topological insulator Cu$_{2}$Bi$_{2}$Se$_{3}$; we find that, up to the first order of momentum, it is impossible to find a proper pairing potential which satisfies the following conditions: (1) it is topologically non-trivial, (2) its energy gap has some point nodes in the $k_{x}k_{y}$-plane, and (3) the Hamiltonian and the corresponding spectral function preserve the threefold rotation symmetry. For two typical pairing potentials, we calculate the corresponding spectral functions and SDOSs. These pairing potentials can act as momentum-dependent analogues to momentum-independent $\Delta_{2}$ and $\Delta_{4}$ proposed in [25] because their spectral functions have the similar shapes. If future experiments, such as the thermal conductivity measurement proposed in [36], do demonstrate that threefold rotation symmetry is broken, more experimental effort is needed to distinguish $\Delta_{4}$ and the momentum-dependent analogue. We also proposed an anisotropic pairing potential which satisfies the three conditions proposed above.

Two things must be clarified. First, the node-contacted bulk band structure is not the unique explanation for ZBPC: a recent paper [37] shows that a fully gapped bulk state with a twisted dispersion of ABS is also possible. Second, the third pairing potential is not the unique choice that satisfies the three conditions proposed above; however, it does serve as the simplest solution. A complete classification of more complicated pairing potentials is far beyond this paper, and it will be given in a future work elsewhere.

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Appendix

To demonstrate that a superconducting topological insulator with pairing potential given by equation (13) is topologically non-trivial, we introduce an intermediate pairing potential

$$\Delta_{c}^{(1)}(k) = \Delta \begin{pmatrix} 0 & A & 0 \\ ZA & 0 & 0 \\ 0 & 0 & -B \end{pmatrix}, \quad (17)$$

where $0 \leq \zeta \leq 1$ is a dimensionless parameter. We divide the integration in equation (12) into two regions: (1) small balls of infinitesimal radius $\epsilon$ around point nodes $k =$

Table 1. The spin and orbital configurations and band gaps of (linear order) momentum-dependent pairing symmetries. Here $\tau_{0}$, $\tau_{1}$, $\tau_{2}$, $\alpha$ are orbital indices, $\tau_{0}$ means pairing electrons belong to the same orbit and the two orbits are anti-symmetric, $\tau_{1}$ means pairing electrons belong to the same orbit and the two orbits are symmetric, $\tau_{2}$ means pairing electrons belong to different orbits.

| Pairing potential | Spin | Orbital | Band gap |
|-------------------|------|---------|----------|
| $\Delta_{1}(k_{x} - ik_{z}) \uparrow \uparrow + i\Delta_{0}(k_{x} + ik_{z}) \downarrow \downarrow + Bk_{z}(\uparrow \downarrow + \downarrow \uparrow)\tau_{0}$ | Triplet | Intra | Gapped |
| $\Delta_{1}(k_{x} - ik_{z}) \uparrow \uparrow + i\Delta_{0}(k_{x} + ik_{z}) \downarrow \downarrow + Bk_{z}(\uparrow \downarrow + \downarrow \uparrow)\tau_{1}$ | Triplet | Intra | Point nodes |
| $\Delta_{1}(k_{x} - ik_{z}) \uparrow \uparrow + i\Delta_{0}(k_{x} + ik_{z}) \downarrow \downarrow + Bk_{z}(\uparrow \downarrow + \downarrow \uparrow)\tau_{2}$ | Triplet | Intra | Gapped |
| $\Delta_{1}(k_{x} - ik_{z}) \uparrow \uparrow + i\Delta_{0}(k_{x} + ik_{z}) \downarrow \downarrow + Bk_{z}(\uparrow \downarrow + \downarrow \uparrow)\tau_{3}$ | Triplet | Inter | Gapped |
| $\Delta_{1}(k_{x} - ik_{z}) \uparrow \uparrow + i\Delta_{0}(k_{x} + ik_{z}) \downarrow \downarrow + Bk_{z}(\uparrow \downarrow + \downarrow \uparrow)\tau_{4}$ | Singlet | Inter | Gapless |
\( \pm (\sqrt{\mu^2 - M^2}/A, 0, 0) \); (2) the remaining regular region. When \( \zeta \) takes values in the interval \((0, 1]\), the bulk gap is not closed and the winding number defined by equation (12) is equal to \((-1)\text{sgn}(\Delta)\) due to the stability of the topology. In this case, there do not exist any singularities, integration over the regular region gives the normal part \((-1)\text{sgn}(\Delta)\) and integration over the two infinitesimal balls gives zero. When \( \zeta \) is equal to zero, integration over the regular region still gives the normal part \((-1)\text{sgn}(\Delta)\) but the integrand is divergent in the infinitesimal balls. Roughly, such a divergence means the winding number (12) is ill defined; however, a careful calculation by choosing \( Q_k \) equal to equation (14) shows that these two singularities are removable singularities, so \( N_w = (-1)\text{sgn}(\Delta) \) still holds in this case.

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] Moore J E 2010 Nature 464 194
[4] König M, Wiedmann S, Brüne C, Roth A, Buhmann H, Molenkamp L W, Qi X L and Zhang S C 2007 Science 318 766
[5] Hsieh D, Qian D, Wray L, Xia Y, Hor Y S, Cava R J and Hasan M Z 2008 Nature 452 970
[6] Hsieh D et al 2009 Nature 460 1101
[7] Xia Y et al 2009 Nature Phys. 5 398
[8] Hsieh D et al 2009 Phys. Rev. Lett. 103 146401
[9] Chen Y L et al 2009 Science 325 178
[10] Fu L, Kane C L and Mele E J 2007 Phys. Rev. Lett. 98 106803
[11] Fu L and Kane C L 2007 Phys. Rev. B 76 045302
[12] Moore J E and Balents L 2007 Phys. Rev. B 75 121306(R)
[13] Qi X L, Hughes T L and Zhang S C 2008 Phys. Rev. B 78 195424
[14] Zhang H, Liu C-X, Qi X L, Dai X, Fang Z and Zhang S C 2009 Nature Phys. 5 438
[15] Liu C-X, Qi X-L, Zhang H, Dai X, Fang Z and Zhang S-C 2010 Phys. Rev. B 82 045122
[16] Schnyder A P, Ryu S, Furusaki A and Ludwig A W W 2008 Phys. Rev. B 78 195125
[17] Kitaev A 2009 AIP Conf. Proc. 1134 22
[18] Wen X-G 2012 Phys. Rev. B 85 085103
[19] Qi X-L, Hughes T L, Raghu S and Zhang S-C 2009 Phys. Rev. Lett. 102 187001
[20] Qi X-L, Hughes T L and Zhang S-C 2010 Phys. Rev. B 81 134508
[21] Fu L and Kane C L 2008 Phys. Rev. Lett. 100 096407
[22] Wilczek F 2009 Nature Phys. 5 614
[23] Hor Y S, Williams A J, Checkelsky J G, Roushan P, Seo J, Xu Q, Zandbergen H W, Yazdani A, Ong N P and Cava R J 2010 Phys. Rev. Lett. 104 057001
[24] Wray L, Xu S-Y, Xia Y, Hor Y, Qian D, Fedorov A, Lin H, Bansil A, Cava R and Hasan M 2010 Nature Phys. 6 855
[25] Sasaki S, Kriener M, Segawa K, Yada K, Tanaka Y, Sato M and Ando Y 2011 Phys. Rev. Lett. 107 217001
[26] Kirzhner T, Lahoud E, Chaska K B, Salman Z and Kanigel A 2012 Phys. Rev. B 86 064517
[27] Yamakage A, Yada K, Sato M and Tanaka Y 2012 Phys. Rev. B 85 180509
[28] Tanaka Y and Kashiwaya S 1995 Phys. Rev. Lett. 74 3451
[29] Kashiwaya S and Tanaka Y 2000 Rep. Prog. Phys. 63 1641
[30] Fu L and Berg E 2010 Phys. Rev. Lett. 105 097001
[31] Sato M 2009 Phys. Rev. B 79 214526
[32] Sato M 2010 Phys. Rev. B 81 220504
[33] Wang Q-H, Wang D and Zhang F-C 2010 Phys. Rev. B 81 035104
[34] Hao L and Lee T K 2011 Phys. Rev. B 83 134516
[35] Sancho M P L, Sancho J M L and Rubio J 1984 J. Phys. F: Met. Phys. 14 1205
[36] Nagai Y, Nakamura H and Machida M 2012 Phys. Rev. B 86 094507
[37] Hsieh T H and Fu L 2012 Phys. Rev. Lett. 108 107005