Majorana Corner Modes and Flat-Band Majorana Edge Modes in Superconductor/Topological-Insulator/Superconductor Junctions

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Recently, superconductors with higher-order topology have stimulated extensive attention and research interest. Higher-order topological superconductors exhibit unconventional bulk-boundary correspondence, thus allowing exotic lower-dimensional boundary modes, such as Majorana corner and hinge modes. However, higher-order topological superconductivity has yet to be found in naturally occurring materials. We investigate higher-order topological superconductivity, which can be achieved by applying magnetic field. When an in-plane Zeeman field is applied to the system, two corner modes appear in the superconducting junction. Furthermore, we also discover a two-dimensional nodal superconducting phase which supports flat-band Majorana edge modes connecting the bulk nodes. Importantly, we demonstrate that zero-energy Majorana corner modes are stable when increasing the thickness of topological insulator thin film.

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The search for topological superconductors which host Majorana zero-energy modes has been one of the central subjects in condensed matter physics,[1–6] since they provide an ideal platform to potential applications in quantum computations based on non-Abelian statistics.[7–14] In 2001, Kitaev proposed to realize Majorana zero-energy modes at the ends of one-dimensional p-wave superconductors.[15] Experimentally, the signatures of zero-energy modes were reported to be observed in spin-orbital coupled semiconductor wires in proximity to s-wave superconductors[16–19] and in the vortex of iron-based superconductors.[20–24] However, conclusive identification of Majorana zero energy modes and scalable fabrication of Majorana networks remain challenging.[5,6,25–28]

Recently, higher-order topological phases of matter, such as the higher-order topological insulators and higher-order topological superconductors, have been identified as a novel topological state, which feature the unconventional bulk-boundary correspondence.[29–51] Generally, the nth higher-order topological superconductors in d dimensions support (d − n)-dimensional gapless boundary excitations with 2 ≤ n ≤ d,[51,52–54] which is in contrast to the d-dimensional conventional topological superconductors with (d − 1)-dimensional boundary excitations.[51,91–93] In this endeavor, Majorana zero-energy modes are supposed to be localized at the corners of a two-dimensional second-order topological superconductor (SOTSC).[51,53,55,58,62,77,80] Due to the fact that naturally occurring topological superconductors are extremely rare, SOTSCs with Majorana Kramers pairs of zero modes or single Majorana zero mode at a corner have been proposed in artificial materials, such as a quantum spin Hall insulators (QSHI) in proximity to a d-wave or an s±-wave superconductor,[53,54] two coupled chiral p-wave superconductors,[55] Rashba spin-orbit coupled s-junction,[56] etc. However, the experimental implementation remains challenging because of the requirements of ideal helical-edge modes of the QSHI, unconventional superconductivities, or complicated junction. Fortunately, topological insulator thin films/s-wave superconductors hybrid structures have been successfully fabricated[51–56] and used to engineer the first-order topological superconductors.[57–59] Even more importantly, a spin-selective Andreev reflection in the vortex of topological insulator/superconductor heterostructure was reported in a scanning tunnel microscope measurement, which is regarded as a fingerprint of Majorana zero-energy modes.[60,61]

In this work, we propose that a topological insulator thin film sandwiched between two s-wave superconductors...
with a phase difference $\pi$ (see Fig. 1) can realize an SOTSC with two localized Majorana corner modes when applying an in-plane magnetic field. We note that the topological insulator-superconductor sandwich structures have been successfully fabricated in the experiment.\cite{102,103,104} To be specific, a topological superconducting phase with gapless helical Majorana edge modes is created when the phase difference of the junction is $\pi$ [see Figs. 2(a) and 2(g)]. Then, the helical Majorana edge modes are gapped out by the in-plane magnetic field due to time-reversal symmetry breaking [see Fig. 2(b)], giving rise to Majorana corner modes [see Figs. 2(e) and 2(h)]. The mechanism of Majorana corners in our proposal is distinguishable from that in Ref. \cite{80}. In Ref. \cite{80}, 1D fermionic helical edge modes of quantum spin Hall insulator are firstly gapped by the s-wave pairing, and then an in-plane magnetic field closes and reopens the edge gap to create Majorana corner modes. Moreover, we find that a nodal topological superconducting phase hosting flat-band Majorana edge modes emerges when tuning the magnetic field. We also show that Majorana corner modes exist when varying the number of layers of topological insulator thin film. Our findings make the superconductor-topological-insulator-superconductor junctions an incredibly fertile platform for exploring topological superconducting phase.

**Model Hamiltonian.** In reciprocal space, the effective Hamiltonian of the superconductor/topological-insulator/superconductor heterostructure can be written as

$$H(k) = -\frac{A}{a} \sin(k_y a) \sigma_y \tau_z + \frac{A}{a} \sin(k_x a) \rho_x \sigma_y \tau_z + M(k) \rho_x \tau_z + \Delta \rho_y \sigma_y \tau_z + V_z \rho_x \sigma_z - \mu \rho_z,$$

in the Nambu basis $C_k = (c_{kl, \uparrow}, c_{kl, \downarrow}, c_{kl, \downarrow}^\dagger, c_{kl, \uparrow}^\dagger)^T$, where $\uparrow$ and $\downarrow$ represent two electron spin directions, and $l = 1, 2$ is the layer index; $\sigma_z, \tau_z$ and $\rho_x(i = x, y, z)$ are the Pauli matrices. In Eq. (1), they act on the spin, layer, and particle-hole spaces, respectively; $\sigma_0, \tau_0$ and $\rho_0$ are the $2 \times 2$ identity matrices. $M(k) = m_0 - \frac{2m_2}{a^2} (2 - \cos(k_x a) - \cos(k_y a))$ describes the mass induced by the hybridization of the top and bottom surfaces of the topological insulator thin film. $A$ is the characteristic parameter of the kinetic energy of the Dirac fermions, and $\mu$ denotes the chemical potential. $\Delta$ is the s-wave pairing amplitude, and the pairing functions of the top and bottom surfaces of the thin film have opposite sign, which makes the setup a Josephson junction with a $\pi$ phase shift. $V_z$ represents the in-plane Zeeman field applied along the $x$ direction. We set the lattice constant $a = 5 \text{ nm}$, $A = 300 \text{ meV} \cdot \text{nm}$ and $m_1 = 150 \text{ meV} \cdot \text{nm}^2$.\cite{105} For our purpose, we set $m_0 = -2 \text{ meV}$ and $\Delta = 2 \text{ meV}$ in this case.

Before turning to discuss the Majorana corner modes in this junction, we would like to give a brief discussion about the case without applied magnetic field. In this case, time-reversal symmetry restores, a topological superconducting phase with gapless helical Majorana edge modes [see Fig. 2(a)] exists when $m_0^2 < \Delta^2 + \mu^2$.\cite{96} The helical edge modes are confirmed by calculating the Bogoliubov quasiparticle energy spectrum and the local density of states for a square nanodisk, as shown in Figs. 2(d) and 2(g).

**Majorana Corner Modes and Flat-Band Edge Modes.** When turning on the in-plane Zeeman field, we find that the gapless helical Majorana edge modes are not stable owing to time-reversal symmetry breaking, in the meantime, an energy gap opens as shown in Fig. 2(b). This gap signals the occurrence of the SOTSC. We compute the Bogoliubov quasiparticle spectrum for a finite-sized sample with a rhombus geometry. The spectrum depicted in Fig. 2(e) shows two degenerate Majorana ingap bound states at zero energy, which reside at the top and bottom corners of the rhombus, respectively, as depicted in Fig. 2(h). These two Majorana corner modes are a smoking-gun signature of the SOTSC.

In the following, we will construct a topological invariant and an edge theory based on the Jackiw–Rebbi mechanism to characterize the zero-energy Majorana corner modes in SOTSC phase.

The Hamiltonian maintains the mirror symmetry: $C_{2x} H(k_x, k_y) C_{2x}^{-1} = H(k_x, -k_y)$ with $C_{2x} = i \sigma_z \tau_z$. Along the mirror invariant axis $k_y = 0$ of the first Brillouin zone (BZ), $H(k_x, k_y = 0)$ commutes with $C_{2x}$ operator. We can use a mirror winding number along this axis to characterize the topological properties of the Majorana corner modes.\cite{76,106,107} The expression of $H(k_x, k_y = 0)$ is

$$H(k_x, 0) = \frac{A}{a} \sin(k_y a) \rho_y \sigma_y \tau_z + M(k_x, 0) \rho_x \tau_z + \Delta \rho_y \sigma_y \tau_z + V_z \rho_x \sigma_z - \mu \rho_z,$$

where $M(k_x, 0) = m_0 - \frac{2m_2}{a^2} (1 - \cos(k_x a))$. $C_{2x}$ has two fourfold degenerate eigenvalues of $\pm 1$. The eigenvectors with eigenvalue of $+1$ are

$$\begin{align*}
\chi_1 &= |\rho_z = 1, \sigma_z = 1, \tau_z = 1\rangle, \\
\chi_2 &= |\rho_z = -1, \sigma_z = 1, \tau_z = 1\rangle, \\
\chi_3 &= |\rho_z = 1, \sigma_z = -1, \tau_z = 1\rangle, \\
\chi_4 &= |\rho_z = -1, \sigma_z = -1, \tau_z = 1\rangle,
\end{align*}
$$

which constitutes the $+1$ eigenspace. The eigenvectors with eigenvalue of $-1$ are

$$\begin{align*}
\chi_5 &= |\rho_z = 1, \sigma_z = 1, \tau_z = -1\rangle, \\
\chi_6 &= |\rho_z = -1, \sigma_z = 1, \tau_z = -1\rangle, \\
\chi_7 &= |\rho_z = 1, \sigma_z = -1, \tau_z = -1\rangle, \\
\chi_8 &= |\rho_z = -1, \sigma_z = -1, \tau_z = -1\rangle,
\end{align*}
$$

which constitutes the $-1$ eigenspace. Projecting $H(k_x, 0)$ into $+1$ eigenspace ($-1$ eigenspace) of $C_{2x}$, we can get a
Hamiltonian in the subspace $H_z(k_x, 0)$ [$H_-(k_x, 0)$]. Using \( \chi_1, \chi_2, \ldots, \chi_8 \) as a new basis set of Hilbert space, we can get $H(k_x, 0) = H_+(k_x, 0) \oplus H_-(k_x, 0)$ with
\[
H_\pm(k_x, 0) = \mp A \kappa \sin(k_x a) \rho_x \sigma_y + M(k_x, 0) \rho_x \sigma_z \pm \Delta \rho_y \sigma_y \pm V \rho_x \sigma_x - \mu \rho_z.
\]
(5)

In the first BZ, we can define the Wilson loop operator $W_{\pm, k_y}$ of $H_\pm(k_x, 0)$, along the mirror invariant axis $k_y = 0$. The mirror winding number $\nu_\pm$ can be written as \cite{106,107}
\[
\nu_\pm = \frac{1}{4 \pi} \log(\det[W_{\pm, k_y}]) \bmod 2.
\]
(6)

When corner modes occur, the mirror winding number $\nu_+ = \nu_- = 1$.\cite{75} The nonzero mirror winding number indicates that the 1D Hamiltonian $H(k_x, 0)$ is topological nontrivial with zero energy modes. Alternatively, we can construct an edge theory to understand corner modes\cite{108} where the Dirac mass domain walls bind zero energy modes. We note that these corner modes are robust when phase shift of the junction is slight deviated from $\pi$, which justify the topological robust of the SOTSC. Furthermore, the corner modes are unaffected by the asymmetry of chemical potential and pairing potential of the top and bottom surfaces. See the Supplementary Materials for details.\cite{108}

Continuing to increase the in-plane Zeeman field $V_z$, we find that the gap closes, and a nodal topological superconducting phase with nodes along the $k_y$-axis is formed. Considering a ribbon geometry with the open boundary condition along $x$ direction, we plot the energy spectrum in Fig. 2(c). Clearly, this nodal phase hosts flat-band Majorana edge modes in between the two bulk nodes. The flat-band Majorana edge modes are confirmed by calculating the energy spectrum and the local density of states of the edge modes, as displayed in Figs. 2(f) and 2(i). We can see that the flat-band Majorana edge modes are located on the top and bottom edges of the sample. The Majorana flat band is featured by a quantized zero-bias conductance peak $2Ne^2/\hbar$ in transport measurement, where $N$ is the number of Majorana zero modes in the Majorana flat band.\cite{109} We note that the Majorana flat band will be tilted,\cite{110} if the paring potential $\Delta$ and chemical potential $\mu$ on the top and bottom surfaces are different. To realize symmetric parameters, we can use highly gate tunable two-dimensional superconductors MoS$_2$ and NbSe$_2$\cite{111,112} to avoid the electrostatic screening by the superconductors.

To capture the topological characteristic of the nodal phase, we further adopt the Wilson loop method\cite{32,33,37,113–116} to calculate the bulk polarization of the system. Since the bulk nodes are located along the $k_y$-axis, by treating $k_y$ as a parameter, the effective Hamiltonian reduces to a one-dimensional Hamiltonian $H_{k_y}(k_x)$. The Wilson loop operator\cite{32,33,37,113–116} along the path in the $k_x$-direction $W_{x,k_y}$ is defined by $W_{x,k_y} = F_{x,k_y+N_x-1}\Delta k_x \cdots F_{x,k_y+1}\Delta k_x F_{x,k_y}$, where $k_y$ is the base point and $N_x$ is the number of unit cells in the $x$-direction. Here, $[F_{x,k_y}]_{mn} = (u_{n,k_y+\Delta k_x}^m u_{m,k_y}^n)$ with the step $\Delta k_x = 2\pi/N_x$, and $|u_{n,k_y}^m\rangle$ represents the occupied Bloch modes.

![Fig. 2. Energy dispersions of nanoribbons geometry along x direction for (a) $V_x = 0$ meV, (b) $V_x = 1.5$ meV, and along y direction for (c) $V_x = 12$ meV. Correspondingly, (d)–(f) the energy spectrum and (g)–(i) local density states, for the sample configurations shown in (g)–(i), where model parameters are the same as those in (a)–(c), respectively. Other parameters are $\mu = 12$ meV and $\Delta = 2$ meV.](image-url)
wave functions with $n = 1, 2, \ldots, N_{\text{occ}}$. $N_{\text{occ}} = N_b/2$ is the number of occupied bands, with $N_b$ the degrees of freedom for each cell. Fixing $k_y$, we can determine the Wilson loop operator $W_{x,k_y}$ on a path along $k_y$. The Wannier center $v_j$ can be determined by the eigenvalues of the Wilson-loop operator $W_{x,k_y}$,

$$W_{x,k_y}|v_j(x,k_y)\rangle = e^{i2\pi v_j}|v_j(x,k_y)\rangle,$$  \hfill (7)

where $j \in \{1, 2, \ldots, N_{\text{occ}}\}$ labels eigenstates $|v_j(x,k_y)\rangle$ as well as components $|v_j(x,k_y)\rangle^n$. Since $k_y$ is fixed, the bulk polarization can be defined as $p = \sum_j v_j \bmod 1$. We plot the bulk polarization as a function of $k_y$ as shown in Fig. 3(b). It is clear that the polarization is quantized to $1/2$ between two nodal points. We remark that the topological characteristic of the nodal phase can be portrayed by the $k_y$-dependent polarization.

Finally, we plot the topological superconducting phase diagram on the plane of $\Delta$ and $V_x$ as shown in Fig. 3(a). The phase boundaries are determined by numerically observing the gap closure of the bulk. The phase boundary between SOTSC and the nodal topological superconductor (NSC) is approximately fitted by the line $V_x = \Delta$. By tuning the in-plane Zeeman field $V_x$, SOTSC phase, NSC phase, and helical topological superconductor (HTSC) phase can be achieved. Figure 3(c) shows the energy spectrum as a function of $V_x$ for the finite-sized sample with rhombus-like geometry in Fig. 2(b).

The effect of thickness on Majorana corner modes. Here, we study how the thickness of the topological insulator thin film affects the Majorana corner modes. In the three-dimensional limit, the bulk Hamiltonian of the intermediate topological insulator can be expressed as

$$H_{3DTI}(k) = \frac{A}{a} \sin(k_x a)\sigma_x \tau_x + \frac{A}{a} \sin(k_y a)\sigma_y \tau_y - \mu$$

$$+ \frac{A}{a} \sin(k_z a)\sigma_z \tau_z + M(k)\tau_z + V_x \sigma_x,$$  \hfill (8)

where $M(k) = m_0 - \frac{2\Delta}{h}[1 - \cos(k_x a)] - \frac{2\Delta}{h}[1 - \cos(k_y a)] - \frac{2\Delta}{h}[1 - \cos(k_z a)]$ with $k = (k_x, k_y, k_z)$, and $V_x$ is the Zeeman field along the $x$-direction.

For simplicity, we consider that the proximity-induced superconducting gap only exits on outermost layers of the top and bottom surfaces of the three-dimensional topological insulator, and the top and bottom superconductors remain a $\pi$ phase shift. In reality, the proximity-induced superconducting potential decreases exponentially and can extend to several layers. However, this will not change the physics discussed here. Considering the confinement of the topological insulator thin film along the $z$-direction, the total Hamiltonian in the Nambu space can be written as

$$H_{3D}(k) = \sum_{z=1}^{N_z-1} b_{k_z,zBound states marked in red dots in (a), (b), and (c), respectively. (g) The gap of bulk spectrum marked in blue dots as a function of number of layers $N_z$. In (a)–(g), we set $a = 5\text{ nm}$, $A = 50\text{ meV}$, $V_x = 1.5\text{ meV}$, $t_x = t_y = 125\text{ meV}\cdot\text{nm}^{-2}$, and $t_z = 100\text{ meV}\cdot\text{nm}^{-2}$.
\[ + \sum_{z=1}^{N_z} b_{k,z}^\dagger (A \sin(k_y a) \sigma_x + A \sin(k_x a) \sigma_z \tau_x \psi_{i,z} + V_s \psi_{i,z} - \mu \psi_{i,z} + \tilde{H}(k) \psi_{i,z} \psi_{i,z}^\dagger + (b_{k,z}^\dagger \Delta \rho_{\sigma_i} \sigma_i b_{k,z} - b_{k,z}^\dagger \Delta \rho_{\sigma_i} \sigma_i b_{k,z}^\dagger) \],

where \( b_{k,z} = [v_{i,z}(k, z), \psi_{i,z}(-k, z)]^T \) with the orbital index \( \ell = (1,2) \) and the spin index \( \alpha = (\uparrow, \downarrow) \). \( \tilde{H}(k) = m_0 - 2t_x (1 - \cos(k_y a)) - 2t_y (1 - \cos(k_x a)) - 2t_z / a^2 \) and \( k_0 = (k_x, k_y) \). Generally, the low energy physics of the topological insulator thin film \( H_{2D}(k) \) can be described by the 2D Hamiltonian \( H(k) \) in Eq. (1). Similar to \( H(k) \), the Hamiltonian \( H_{2D}(k) \) preserves the mirror-like symmetry: \( C_{2x} H_{2D}(k_x, k_y) C_{2x}^{-1} = H_{2D}(k_y, -k_x) \) with \( C_{2x} = i \sigma_x \mathcal{U} \). Here \( N_x \times N_y \) anti-diagonal matrix \( \mathcal{U} = \delta_{k_x, \pi} \mathcal{U}_1 \) is defined in real space with Kronecker delta \( \delta_{k_x, \pi} \). Therefore, it is natural to expect that the Majorana corner modes will be sustainable with the varying thickness. For \( \Delta \neq 0 \), \( V_s = 0 \), the system is a helical topological superconductor, if the Fermi energy is outside the surface gap. By turning on the Zeeman field \( V_s > 0 \), the gapless Majorana corner modes are observed, as shown in Fig. 4. Figures 4(a), 4(b), and 4(c) show the energy spectrum (only the 20 eigenenergies with the smallest absolute value are shown) of \( N_x \times N_y = 200 \times 200 \) rhombus-shaped sample. Again, two zero-energy ingap Majorana bound states appear (marked by the red dots in Fig. 4). Figures 4(d), 4(e), and 4(f) are the local density of the two zero-energy Majorana bound states in Figs. 4(a), 4(b), and 4(c), respectively. We can see that the two Majorana bound states are also localized at two opposite corners in the \( xy \)-plane, but extended in hinge of the side surface along the \( z \)-direction. In order to explore the relationship between the gap of side-surface spectrum [marked in blue dots in Figs. 4(a)–4(c)] and the number of layers \( N_z \), we plot the band gap for different numbers of layers \( N_z \) in Fig. 4(g). The band gap decreases rapidly with the increase of \( N_z \) due to the decreasing finite-size confinement along the \( z \)-direction.

In conclusion, we have illustrated that an SOTSC with two Majorana corner modes is realized in topological insulator thin film based superconducting junctions with a \( \pi \) phase shift when an in-plane Zeeman field is applied. We employ the mirror winding number to characterize the second-order topology of Majorana corner modes. We also analytically deduce an edge theory for the Majorana corner modes by using the perturbation theory. By tuning the Zeeman field, we also observe a nodal superconducting phase hosting flat-band Majorana edge modes, whose bulk topology can be captured by a \( k \)-dependent polarization. Lastly, we demonstrate how the thickness of topological insulator thin films affects the Majorana corner modes and their spatial distribution.

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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] Alicea J 2012 Rep. Progress Physica 75 076501
[4] Beenakker C 2013 Annu. Rev. Condens. Matter Phys. 4 113
[5] Lutchyn R M, Bakkers E P, Kouwenhoven L P, Krogstrup P, Marcus C M, and Oreg Y 2018 Nat. Rev. Mater. 3 52
[6] Flensberg, k, von Oppen F, and Stern A 2021 Nat. Rev. Mater. 6 944
[7] Ivanov D A 2001 Phys. Rev. Lett. 86 268
[8] Kitaev A Y 2001 Ann. Phys. 303 32
[9] Nayak C, Simon S H, Stern A, Freedman M, and Das S S 2008 Rev. Mod. Phys. 80 1083
[10] Alicea J, Oreg Y, Refael G, von Oppen F, and Fisher M P A 2011 Nat. Phys. 7 412
[11] Flensberg K 2011 Phys. Rev. Lett. 106 090503
[12] van Heck B, Akhmerov A R, Hassler F, Burrello M, and Beenakker C W J 2012 New J. Phys. 14 035019
[13] Aasen D, Hell M, Mishmash R V, Higginbotham A, Danon J, Leijnse M, Jespersen T S, Folk J A, Marcus C M, Flensberg K, and Alicea J 2016 Phys. Rev. Lett. 116 067001
[14] Zhao L B, Wang Z D, Shen R, Sheng L, Wang B G, and Xing D Y 2017 Chin. Phys. Lett. 34 057401
[15] Kitaev A Y 2001 Phys. Usp. 44 131
[16] Mourik V, Zuo K, Frolov S M, Plissard S R, Bakkers E P A M, and Kouwenhoven L P 2012 Science 336 1003
[17] Das A, Ronen Y, Most Y, Oreg Y, Heiblum M, and Shtrikman H 2012 Nat. Phys. 8 887
[18] Deng M T, Yu C L, Huang G Y, Larsson M, Caroff P, and Xu H Q 2012 Nano Lett. 12 6414
[19] Lutchyn R M, Sau J D, and Das S S 2010 Phys. Rev. Lett. 105 077001
[20] Zhang P, Yaji K, Hashimoto T, Ota Y, Kondo T, Ozaaki K, Wang Z, Wen Z, Gu G D, Ding H, and Shin S 2018 Science 360 182
[21] Kong L Y, Zhu S Y, Papaj M, Chen H, Cao L, Isobe H, Xing Y, Liu W, Wang D, Fan P, Sun Y, Du S, Schneeloch J, Zhang R, Gu G, Fu L, Gao H J, and Ding H 2019 Nat. Phys. 15 1181
[22] Wang D F, Kong L Y, Fan P, Chen H, Zhu S Y, Liu W Y, Cao L, Sun Y J, Du S X, Schneeloch J, Zhang R D, Gu G D, Fu L, Ding H, and Gao H J 2018 Science 362 333
[23] Liu Q, Chen C, Zhang T, Peng R, Yan Y J, Wen C H P, Lou X, Huang Y L, Tian J P, Dong X L, Wang G W, Bao W C, Wang Q H, Yin Z F, Zhao Z X, and Feng D L 2018 Phys. Rev. X 8 041056
[24] Zhu S Y, Kong L Y, Cao L, Chen H, Papaj M, Du S, Xing Y Q, Liu W Y, Wang D F, Shen C M et al. 2020 Science 367 189
[25] Karzig T, Knapp C, Lutchyn R M, Bonderson P, Hastings M B, Nayak C, Alicea J, Flensberg K, Plugge S, Oreg Y, Marcus C M, and Freedman M H 2017 Phys. Rev. B 95 235305
[26] Chen C, Liu Q, Zhang T Z, Li D, Shen P P, Dong X L, Zhao Z X, Zhang T, and Feng D L 2010 Chin. Phys. Lett. 27 057043
[27] Liu S, Nie S M, Qi Y P, Guo Y F, Yuan H T, Yang L X, Chen Y L, Wang M X, and Liu Z K 2021 Chin. Phys. Lett. 38 077302
[28] Pan D, Song H, Zhang S, Liu L, Wen L, Liao D, Zhuo R, Wang Z, Zhang Z, Yang S, Ying J, Miao W, Shang R, Zhang H, and Zhao J 2022 Chin. Phys. Lett. 39 058101
[29] Peng Y, Yao B, and von Oppen F 2017 Phys. Rev. B 95 235143
[30] Song Z D, Fang Z, and Fang C 2017 Phys. Rev. Lett. 119 246402
[31] Langbehn J, Peng Y, Trifunovic L, von Oppen F, and Brouwer P W 2017 Phys. Rev. Lett. 119 246401
[32] Benalcazar W A, Bernevig B A, and Hughes T L 2017 Phys. Rev. B 96 245115
[33] Benalcazar W A, Bernevig B A, and Hughes T L 2017 Science 357 61
