Aharonov-Bohm effect in three-dimensional higher-order topological insulators

Kun Luo,1 Hao Geng,1,2 Li Sheng,1,2 Wei Chen,1,2 and D. Y. Xing1,2

1 National Laboratory of Solid State Microstructures and school of Physics, Nanjing University, Nanjing, 210093, China
2 Collaborative Innovation Center of Advanced Microstructures, Nanjing University, Nanjing 210093, China

(Dated: August 27, 2021)

The 1D hinge states are the hallmark of the 3D higher-order topological insulators (HOTI), which may lead to interesting transport properties. Here, we study the Aharonov-Bohm (AB) effect in the interferometer constructed by the hinge states in the normal metal-HOTI junctions with a transverse magnetic field. We show that the AB oscillation of the conductance can clearly manifest the spatial configurations of such hinge states. The magnetic fluxes encircled by various interfering loops are composed of two basic ones, so that the oscillation of the conductance by varying the magnetic field contains different frequency components universally related to each other. Specifically, the four dominant frequencies \( \omega_{x,y} \) and \( \omega_{\pm} \) satisfy the relations \( \omega_{\pm} = \omega_x \pm \omega_y \), which generally holds for different magnetic field, sample size, bias voltage and weak disorder. Our results provide a unique and robust signature of the hinge states and pave the way for exploring AB effect in the 3D HOTI.

I. INTRODUCTION

Over the past two decades, topological phases of matter such as topological insulator and superconductor have become an active research field of condensed matter physics [1, 2]. These materials are characterized by the nontrivial band topology and the resultant gapless (d-1)-dimensional edge states. Very recently, the concepts of higher-order topological insulators (HOTI) and superconductors are theoretically proposed, which are featured by the (d-2)-dimensional edge states [3, 17]. Specifically, for the 3D HOTI there exist 1D gapless states along the hinges of the sample, so-called hinge states, while the surface and the bulk states are both insulating. Recent progresses have shown the evidences of the hinge states in bismuth by the scanning-tunnelling spectroscopy and Josephson interferometry [17], which pave the way for exploring more intriguing properties of such topological states in the HOTI.

The 1D nature of the hinge states indicates that it is a good playground for exploring various interference effects, such as Aharonov-Bohm (AB) and Fabry-Pérot interferometers [18–23]. Actually, the chiral edge states of the quantum Hall phase have become an important platform for the study of mesoscopic physics, in which a variety of novel phenomena have been observed [20, 24, 25] due to its long coherence length and high adjustability. Compared with the chiral edge states, the hinge states in the HOTI open additional possibilities for the implementation of novel effects due to their 3D configurations, which enrich the way of interfering in real space. Moreover, such effects cannot be realized in any 2D systems, which in turn, can serve as the deterministic evidence of the hinge states.

The manifestation of the AB effect in an electron system is the periodic oscillation of conductance as the closed trajectory of electrons encircles a magnetic flux \( \Phi \) [27, 32]. The dominant period of oscillation is equal to the flux quantum \( \Phi_0 = h/e \), with a main frequency \( 2\pi/\Phi_0 \). The frequency of oscillation can be found by taking the fast Fourier transform (FFT) of the conductance pattern [28, 30]. Recently, AB effect has been used as an effective way to detect edge states in various topological systems, such as edge states of topological insulators [33–36], Majorana fermions of topological superconductors [37, 40], surface states of topological semimetal [38, 42], and non-Abelian anyons of fractional quantum Hall systems [23, 43, 46].

In this work, we investigate the AB effect in the interferometer composed of the hinge states of the quadrangular HOTI by imposing an external magnetic field. The insulating bulk and surface states indicate that the electron can only propagate along the hinges of the sample, by which the enclosed magnetic flux can lead to a coherent oscillation of the transmission probability. Different from the edge states in any 2D systems, the 3D network of the hinge states results in peculiar interfering trajectories, which relies not only on the magnitude of
the magnetic field but also on its orientation. The AB interferometer is sketched in Fig. 1(a), where the HOTI is connected to two leads made of the normal metal and a magnetic field $B = (B_x, B_y, B_z) = B(\cos \theta, \sin \theta)$ is applied in the $x$-$y$ plane with $\theta$ being the polar angle. The electrons injected from the leads propagate along four chiral hinge states, which comprise a variety of interfering loops; see Figs. 1(b) and 1(c). The elemental interfering loops shown in Figs. 1(b) and 1(c) are exactly the boundaries of the $(\pm 1,0,0)$ and $(0,\pm 1,0)$ surfaces. The basic loops in Figs. 1(b) and 1(c) encircle a magnetic flux of $\Phi$ naturally appear in the oscillating pattern of the conductance as the magnetic field $B$ varies. Interestingly, the magnetic flux in other interfering loops can all be interpreted by the two elemental ones, among which two typical loops in Figs. 1(d) and 1(e) contain a flux of $\Phi_x, y = \Phi_x \pm \Phi_y$, and the corresponding oscillating frequency components satisfy $\omega_{\pm y} = \omega_x \pm \omega_y$. It turns out that the aforementioned four interfering loops and the corresponding oscillating frequencies dominate the coherent oscillation of the conductance. The relation $\omega_{x+y} = \omega_x \pm \omega_y$ generally holds independent of various parameters such as the magnetic field, sample size and the energy of electron, thus providing a universal and deterministic signature of the hinge states and HOTI.

The rest of this paper is organized as follows. In Sec. II, we elucidate the model of the HOTI adopted in our work. In Sec. III, we apply the scattering matrix approach to analyze the coherent transport through the interferometer and the AB oscillation of the conductance. Detailed numerical simulations on the lattice model are conducted in Sec. IV, which verify the universality of the physical results. Finally, a brief summary and outlook are given in Sec. V.

II. MODEL OF HOTI

We adopt the model of 3D chiral HOTI introduced by Schindler et al. [7] as

$$H_{\text{HOTI}} = \left( M + t \sum_i \cos k_i \right) \tau_z \sigma_0 + \Delta_1 \sum_i \sin k_i \tau_z \sigma_i + \Delta_2 (\cos k_x - \cos k_y) \tau_y \sigma_0,$$

where $\sigma_i = x, y, z$ and $\tau_i$ are the Pauli matrices acting on the spin and orbital space, respectively. For $1 < |M/t| < 3$ and $\Delta_1, \Delta_2 \neq 0$, the system lies in a chiral 3D HOTI phase. The energy spectra are gapped in both the bulk and four surfaces parallel to the $z$-axis. Importantly, the mass term is opposite in sign between adjacent surfaces that results in the Jackiw-Rebbi-type bound states [17] propagating only along the $\pm z$-direction, or the so-called topological hinge states. Time-reversal symmetry is broken in Eq. (1), so that the hinge states are unidirectional or chiral, without any backscattering states within a given hinge. Notably, gapless Dirac cones protected by the $C_{\text{d}}^2 T$ symmetry persist on the surfaces perpendicular to the $z$-axis [7]. Therefore, it is beneficial to explore pure signature of the hinge states through the transport in the $z$-direction.

III. SCATTERING MATRIX ANALYSIS

In this section, we study the coherent transport of electrons through the interferometer sketched in Fig. 1(a) based on the low-energy effective model of the hinge states using the scattering matrix approach. The scattering matrix of the whole interferometer can be obtained by combining those at two normal metal-HOTI interfaces and the matrix of phase accumulation during propagation in the hinge states. The matrix at the lower interface [cf. Fig. 1(a)] can be parameterized as

$$S_l = \begin{pmatrix} r_1 & r_3 & t_1' & t_3' \\ r_2 & r_4 & t_2' & t_4' \\ t_1 & t_3 & r_1' & r_3' \\ t_2 & t_4 & r_2' & r_4' \end{pmatrix},$$

(2)

which relates the incoming ($a_l$) and outgoing ($b_l$) waves in the normal metal and the HOTI via $b_l = S_l a_l$. The matrix is assumed to be $4 \times 4$ such that two incoming/outgoing waves are taken into account on both sides. For the HOTI, the number of channels corresponds to that of pairs of the hinge states. The unitary condition $S_l S_l^\dagger = 1$ is ensured by the law of current conservation. Here, $r_{1,\ldots,4}$ and $t_{1,\ldots,4}$ are the transmission amplitudes from the normal metal to the chiral hinge states of the HOTI and $r_{1,\ldots,4}$ are the corresponding reflection amplitudes. The scattering amplitudes corresponding to the incident waves from the hinge states of HOTI are defined by $t_{1,\ldots,4}, r_{1,\ldots,4}$ in a similar way. The scattering matrix for the upper interface can be defined as

$$S_u = \begin{pmatrix} r_1^u & r_3^u & t_1'^u & t_3'^u \\ r_2^u & r_4^u & t_2'^u & t_4'^u \\ t_1^u & t_3^u & r_1'^u & r_3'^u \\ t_2^u & t_4^u & r_2'^u & r_4'^u \end{pmatrix},$$

(3)

The phase modulation of the wave function due to the magnetic field can be described by the matrix as

$$S_m = \begin{pmatrix} 0 & 0 & e^{i\phi_1} & 0 \\ 0 & 0 & 0 & e^{-i\phi_2} \\ e^{-i\phi_2} & 0 & 0 & 0 \\ 0 & e^{i\phi_1} & 0 & 0 \end{pmatrix},$$

(4)

where the phases $\phi_1, \phi_2, \tilde{\phi}_1, \tilde{\phi}_2$ are related by the magnetic fluxes through $\phi_1 + \phi_2 = \phi_2 + \tilde{\phi}_2 = \phi_x = 2\pi \Phi_x/\Phi_0, \phi_1 - \phi_2 = \phi_2 - \phi_1 = \phi_y = 2\pi \Phi_y/\Phi_0, \phi_1 + \phi_2 = \phi_z = 2\pi \Phi_z/\Phi_0$. The scattering matrix for the whole interferometer sketched in Fig. 1(a) can be parameterized as

$$S = S_u S_l^\dagger S_m S_l.$$
\[ \phi_{x+y} = 2\pi \Phi_{x+y}/\Phi_0 \text{ and } \phi_1 + \phi_2 = \phi_{x-y} = 2\pi \Phi_{x-y}/\Phi_0 \] with \( \phi_{x,y} \) and \( \phi_{x\pm y} \) being gauge invariant.

By combining three matrices \( S_l, S_m, S_u \) in a standard way we obtain the total scattering matrix for the whole system. Here, we focus on the periods of the AB oscillation and an overall phase shift of the pattern is unimportant. Therefore, we can choose \( S_l, S_u \) to be real for simplicity which will not change the main results. For an electron incident from the lower terminal, its transmission probability \( T \) to upper terminal is obtained after some algebra as

\[
T = F^{-1} \left[ C + C_X \cos \phi_x + C_Y \cos \phi_y + C'_{XY} \cos \phi_{x+y} + C''_{XY} \cos \phi_{x-y} \right],
\]

\[
F = M_0 + M_{XY} \cos \phi_{x+y} + M'_{XY} \cos \phi_{x-y} + M_{2X} \cos (2\phi_x) + M_{2Y} \cos (2\phi_y) - M_X \cos \phi_x - M_Y \cos \phi_y,
\]

where the explicit forms of the relevant parameters are given in Appendix A. The numerator of the transmission in Eq. (5) shows that there are four dominant periodic terms contributed by four interference loops in Figs. 1(b)-(e) which correspond to four frequencies related by \( \omega_{x,y}, \omega_{x\pm y} = \omega_x \pm \omega_y \). Note that such relations are stabilized by the spatial configurations of the hinge states, thus offer a clear and robust signature for its detection, which relies little on the sample details and the energy. Although magnetic field in different directions will change the values of frequencies it does not affect the general relations between them.

Next, we provide numerical verification of such an observation using specific scattering amplitudes. The AB oscillation of the conductance as a function of \( B \) and its FFT spectrum are shown in Figs. 2(a) and 2(b), respectively. The polar angle of the magnetic field is set to \( \theta = \pi/6 \) and the unit of frequency is chosen as \( 1/B_0 \) with \( B_0 = \Phi_0/(2\pi S) \) and \( S = S_x = S_y \) the surface area. One can find multiple periods in the conductance pattern in Fig. 2(a). The FFT spectrum in Fig. 2(b) shows that there are four dominant frequencies with \( \omega_x = 0.86/B_0, \omega_y = 0.5/B_0, \omega_{x-y} = 0.36/B_0, \omega_{x+y} = 1.36/B_0 \), which conforms the aforementioned relation \( \omega_{x\pm y} = \omega_x \pm \omega_y \). Higher frequencies such as \( \omega_{2x}, \omega_{2y} \) should also appear as that in the conventional 2D AB effect. In stark contrast, the frequencies \( \omega_{x\pm y} \) can only exist in the 3D HOTI, which thus provides a unique evidence of the hinge states.

IV. LATTICE MODEL SIMULATION

Based on the scattering matrix analysis, we see that there are four dominant frequencies satisfying universal relations \( \omega_{x\pm y} = \omega_x \pm \omega_y \). In this section, we perform numerical simulation to give rigorous results. We write

\[
H_{\text{HOTI}}^{\text{Lattice}} = \sum_i c_i^\dagger M \sigma_0 \tau_z c_i + \left\{ \sum_i c_i^\dagger e^{i\varphi_x} \frac{\tau_y}{2} (\Delta_2 \sigma_0 \tau_z + t \sigma_0 \tau_x + i \Delta_1 \sigma_2 \tau_x) c_i 
- \sum_i c_i^\dagger (\Delta_2 \sigma_0 \tau_y + t \sigma_0 \tau_x + i \Delta_1 \sigma_2 \tau_x) c_i 
+ \sum_i c_i^\dagger e^{i\varphi_y} \frac{\tau_x}{2} (\sigma_0 \tau_z + i \Delta_1 \sigma_x \tau_z) c_i + h.c. \right\},
\]

where \( c_i = (c_{0,\uparrow,i}, c_{0,\downarrow,i}, c_{a,\uparrow,i}, c_{a,\downarrow,i}) \) are the annihilate operators at lattice site \( i \) with two spin \( (\uparrow, \downarrow) \) and two orbit \( (a, b) \) components. The Peierls phase \( \varphi_{x,z} = \int_{r_1}^{r_2} A(r) \cdot dr \), where \( A(r) = (B_y r_z, 0, B_x r_y) \) is the vector potential under Landau gauge. The lattice model of the normal...
and polar angle conductance oscillation for different incident energy (\(\theta\)). For different incident energies, the oscillation patterns in Fig. 3(a) look like those in Figs. 2(b). For different spectral data, the dominant frequencies, we perform FFT calculation whose results are consistent with those by the scattering matrix analysis in Figs. 2(b).

The lattice of the HOTI is set to a 30° lattice. The magnetic field \(B\) is calculated using KWANT package and the AB effect should be robust, the same as that in the quantum Hall edge states. We show numerical results for the disorder distributed in the whole HOTI region and is parallel to the \(x-y\) plane.

In Fig. 4, we present more general results by varying both the polar angle \(\theta\) and the thickness of the HOTI in the \(z\)-direction. Each pair of parameters generate one point in both Figs. 4(a) and 4(b), with its coordinates extracted in the same way as done in Fig. 3(b). The reference planes therein correspond to the frequency rule \(\omega_x = \omega_z\) for \(\omega_y\). One can see that the numerical results labeled by the black and orange dots are well located around the reference planes, which indicates the universality of the frequency rule. Note that there are a few dots of negative frequencies in Fig. 4(b) for \(\omega_y\). In experiments, one should rather measure \(|\omega_x - \omega_y|\) instead.

Disorder generally exists in real samples and it can be expected that topological chiral hinge states and thus the AB effect should be robust, the same as that in the quantum Hall edge states. We show numerical results for the disorder distributed in the whole HOTI region in Fig. 5 with different strength. For weak disor-
der strength $\alpha < 0.7|t|$ (the gap of the bulk states), the oscillation pattern and frequency rules $\omega_{x\pm y} = \omega_x \pm \omega_y$ retain. For strong disorder $\alpha > 0.7|t|$, the oscillation pattern quenches stemming from disorder induced coupling between the surface/bulk states and the hinge states. Therefore, as long as the disorder in the sample of HOTI is not too strong, the AB effect can be hopefully observed. Similar conclusion also holds for the surface roughness. Although the AB effect is quite robust against the disorder effect, the observation should be carried out within the phase coherence length of the system. The dephasing effect always reduces the visibility of the coherent oscillation until it vanishes [30,50]. One more remark is that the interference here is all of the AB type without Al'tshuler-Aronov-Spivak (AAS) type contribution [51,53]. The model in Eq. (1) breaks time-reversal symmetry so that the AAS effect is absent.

V. SUMMARY AND OUTLOOK

In summary, we have investigated the AB effect in the chiral hinge states of the 3D HOTI. Due to the spatial configurations of the hinge states, new types of interfering loops appear compared with the 2D AB interference. Importantly, we predict a universal relationship $\omega_{x\pm y} = \omega_x \pm \omega_y$ between the dominant oscillating frequencies, which offers a unique signature of the hinge states as well as the HOTI. Our study can be generalized straightforwardly to AB effect in the 3D HOTI with helical hinge states.

Note added. Recently, we became aware of a related work [54], which focuses on different aspects.

ACKNOWLEDGMENTS

This work was supported by the National Natural Science Foundation of China under Grant No. 12074172 (W.C.), No. 11674160 and No. 11974168 (L.S.), the startup grant at Nanjing University (W.C.), the State Key Program for Basic Researches of China under Grants No. 2017YFA0303203 (D.Y.X.) and the Excellent Programme at Nanjing University.

Appendix A: specific forms of coefficient for analysis calculation

In Eq. (5) of the main text, the coefficients are expressed as

$$C = (t_1^2 + t_2^2)(W_0^2 + W_1^2 + W_2^2 + X_0^2 + X_1^2 + X_2^2) + (t_3^2 + t_4^2)(Y_0^2 + Y_1^2 + Y_2^2 + Z_0^2 + Z_1^2 + Z_2^2),$$

$$C_X = 2(W_0W_1 + X_0X_1)(t_1^2 + t_2^2) + 2(Y_0Y_1 + Z_0Z_1)(t_3^2 + t_4^2) + 2(W_0Y_2 + W_1Y_3 + X_0Z_2 + X_1Z_3)(t_1t_2 + t_3t_4),$$

$$C_Y = 2(W_0W_2 + X_0X_2)(t_1^2 + t_3^2) + 2(Y_0Y_2 + Z_0Z_2)(t_2^2 + t_4^2) + 2(W_0Y_1 + W_1Y_3 + X_0Z_2 + X_1Z_3)(t_1t_2 + t_3t_4),$$

$$C_{XY} = 2(W_0W_3 + X_0X_3)(t_1t_2 + t_3t_4),$$

$$C'_{XY} = 2(W_1W_2 + X_1X_2)(t_1^2 + t_3^2) + 2(Y_1Y_2 + Z_1Z_2)(t_2^2 + t_4^2) + 2(W_1Y_3 + W_2Y_4 + X_1Z_2 + X_2Z_3)(t_1t_2 + t_3t_4),$$

$$M_{XY} = 2M_1M_4 + 2M_2M_3,$$

$$M'_{XY} = 2M_1M_3 + 2M_2M_4,$$

$$M_2 = 2M_1M_2,$$

$$M_4 = 2M_1M_4,$$

$$M_C = M_0^2 + M_1^2 + M_2^2 + M_3^2 + M_4^2,$$

(A1)

which contain the parameters defined by the elements of the scattering matrices as

$$W_0 = t_1^u, W_1 = t_3^u + t_1^v - t_1^u r_3^{r_1}, W_2 = t_3^u + t_1^v - t_1^u r_3^{r_4},$$

$$X_0 = t_2^u, X_1 = t_4^u + t_2^v - t_2^u r_2^{r_1}, X_2 = t_4^u + t_2^v - t_2^u r_2^{r_4},$$

$$Y_0 = t_3^u, Y_1 = t_5^u + t_3^v - t_3^u r_3^{r_1}, Y_2 = t_5^u + t_3^v - t_3^u r_3^{r_4},$$

$$Z_0 = t_4^u, Z_1 = t_6^u + t_4^v - t_4^u r_4^{r_1}, Z_2 = t_6^u + t_4^v - t_4^u r_4^{r_4},$$

$$M_0 = 1 + r_1^{r_1} r_1^{r_1} r_1^{r_4} + r_4^{r_4} r_4^{r_4} - r_2^{r_2} r_2^{r_2} - r_3^{r_3} r_3^{r_3} - r_1^{r_1} r_1^{r_4} - r_3^{r_3} r_3^{r_4},$$

(A2)

$$M_1 = r_2^{r_3}, M_2 = r_3^{r_2}, M_3 = r_4^{r_4}, M_4 = r_1^{r_1}.$$
[1] X.-L. Qi and S.-C. Zhang, Rev. Mod. Phys. 83, 1057 (2011).
[2] M. Z. Hasan and C. L. Kane, Rev. Mod. Phys. 82, 3045 (2010).
[3] W. A. Benalcazar, B. A. Bernevig, and T. L. Hughes, Science 357, 61 (2017).
[4] W. A. Benalcazar, B. A. Bernevig, and T. L. Hughes, Phys. Rev. B 96, 245115 (2017).
[5] Z. Song, Z. Fang, and C. Fang, Phys. Rev. Lett. 119, 246402 (2017).
[6] J. Langbehn, Y. Peng, L. Trifunovic, F. von Oppen, and P. W. Brouwer, Phys. Rev. Lett. 119, 246401 (2017).
[7] F. Schindler, A. M. Cook, M. G. Vergniory, Z. Wang, S. S. Parkin, B. A. Bernevig, and T. Neupert, Science advances 4, eaat0346 (2018).
[8] S. Franca, J. van den Brink, and I. C. Fulga, Phys. Rev. B 98, 201114 (2018).
[9] M. J. Park, Y. Kim, G. Y. Cho, and S. Lee, Phys. Rev. Lett. 123, 216803 (2019).
[10] E. Khalat, Phys. Rev. B 97, 205136 (2018).
[11] D. Yu, R.-X. Zhang, and S. Das Sarma, Phys. Rev. Research 2, 043223 (2020).
[12] R.-X. Zhang, Y.-T. Hsu, and S. Das Sarma, Phys. Rev. B 102, 094503 (2020).
[13] Z. Yan, Phys. Rev. Lett. 123, 177001 (2019).
[14] D. Călugăru, V. Juričić, and B. Roy, Phys. Rev. B 99, 041301 (2019).
[15] Z. Yan, Phys. Rev. B 100, 205406 (2019).
[16] Z. Yan, F. Song, and Z. Wang, Phys. Rev. Lett. 121, 096803 (2018).
[17] F. Schindler, Z. Wang, M. G. Vergniory, A. M. Cook, A. Murani, S. Sengupta, A. Y. Kasumov, R. Debloch, S. Jeon, I. Drozdov, et al., Nature physics 14, 918 (2018).
[18] W. Liang, M. Bockrath, D. Bozovic, J. H. Hafner, M. Tinkham, and H. Park, Nature 411, 665 (2001).
[19] B. J. van Wees, L. P. Kouwenhoven, C. J. P. M. Harmans, J. G. Williamson, C. E. Timmering, M. E. I. Broekaart, C. T. Foxon, and J. J. Harris, Phys. Rev. Lett. 62, 2523 (1989).
[20] Y. Ji, Y. Chung, D. Sprinzak, M. Heiblum, D. Mahalu, and H. Shtrikman, Nature 422, 415 (2003).
[21] N. Ofek, A. Bid, M. Heiblum, A. Stern, V. Umansky, and D. Mahalu, Proceedings of the National Academy of Sciences 107, 5276 (2010).
[22] D. T. McClure, Y. Zhang, B. Rosenow, E. M. Levenson-Falk, C. M. Marcus, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. 103, 206806 (2009).
[23] J. Nakamura, S. Fällåh, H. Sahasrabudhe, R. Rahman, S. Liang, G. C. Gardner, and M. J. Manfra, Nature Physics 15, 563 (2019).
[24] M. Henry, S. Oberholzer, C. Strunk, T. Heinzel, K. Esslin, M. Holland, and C. Schönenberger, Science 284, 296 (1999).
[25] I. Neder, N. Ofek, Y. Chung, M. Heiblum, D. Mahalu, and V. Umansky, Nature 448, 333 (2007).
[26] E. Weisz, H. Choi, I. Sivan, M. Heiblum, Y. Gefen, D. Mahalu, and V. Umansky, Science 344, 1363 (2014).
[27] Y. Aharonov and D. Bohm, Phys. Rev. 115, 485 (1959).
[28] R. A. Webb, S. Washburn, C. P. Umbach, and R. B. Laibowitz, Phys. Rev. Lett. 54, 2696 (1985).
[29] G. W. Holloway, D. Shiri, C. M. Haapanaki, K. Willick, G. Watson, R. R. LaPierre, and J. Baugh, Phys. Rev. B 91, 045422 (2015).
[30] A. Lahiri, K. Gharavi, J. Baugh, and M. Muralidharan, Phys. Rev. B 98, 125417 (2018).
[31] I. L. Aleiner, A. V. Andreev, and V. Vinokur, Phys. Rev. Lett. 114, 076802 (2015).
[32] Y. Tserkovnyak and B. I. Halperin, Phys. Rev. B 74, 245327 (2006).
[33] H. Peng, K. Lai, D. Kong, S. Meister, Y. Chen, X.-L. Qi, S.-C. Zhang, Z.-X. Shen, and Y. Cui, Nature materials 9, 225 (2010).
[34] J. H. Bardarson, P. W. Brouwer, and J. E. Moore, Phys. Rev. Lett. 105, 156803 (2010).
[35] Y. Zhang and A. Vishwanath, Phys. Rev. Lett. 105, 206601 (2010).
[36] F. Schindler, A. M. Cook, M. G. Vergniory, Z. Wang, S. S. Parkin, B. A. Bernevig, and T. Neupert, Science advances 4, eaat0346 (2018).
[37] S. Franca, J. van den Brink, and I. C. Fulga, Phys. Rev. B 98, 201114 (2018).
[38] M. J. Park, Y. Kim, G. Y. Cho, and S. Lee, Phys. Rev. Lett. 123, 216803 (2019).
[39] E. Khalat, Phys. Rev. B 97, 205136 (2018).
[40] D. Yu, R.-X. Zhang, and S. Das Sarma, Phys. Rev. Research 2, 043223 (2020).
[41] R.-X. Zhang, Y.-T. Hsu, and S. Das Sarma, Phys. Rev. B 102, 094503 (2020).
[42] Z. Yan, Phys. Rev. Lett. 123, 177001 (2019).
[43] D. Călugăru, V. Juričić, and B. Roy, Phys. Rev. B 99, 041301 (2019).
[44] Z. Yan, Phys. Rev. B 100, 205406 (2019).
[45] Z. Yan, F. Song, and Z. Wang, Phys. Rev. Lett. 121, 096803 (2018).
[46] F. Schindler, Z. Wang, M. G. Vergniory, A. M. Cook, A. Murani, S. Sengupta, A. Y. Kasumov, R. Debloch, S. Jeon, I. Drozdov, et al., Nature physics 14, 918 (2018).
[47] W. Liang, M. Bockrath, D. Bozovic, J. H. Hafner, M. Tinkham, and H. Park, Nature 411, 665 (2001).
[48] B. J. van Wees, L. P. Kouwenhoven, C. J. P. M. Harmans, J. G. Williamson, C. E. Timmering, M. E. I. Broekaart, C. T. Foxon, and J. J. Harris, Phys. Rev. Lett. 62, 2523 (1989).
[49] Y. Ji, Y. Chung, D. Sprinzak, M. Heiblum, D. Mahalu, and H. Shtrikman, Nature 422, 415 (2003).
[50] N. Ofek, A. Bid, M. Heiblum, A. Stern, V. Umansky, and D. Mahalu, Proceedings of the National Academy of Sciences 107, 5276 (2010).
[51] D. T. McClure, Y. Zhang, B. Rosenow, E. M. Levenson-Falk, C. M. Marcus, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. 103, 206806 (2009).
[52] J. Nakamura, S. Fällåh, H. Sahasrabudhe, R. Rahman, S. Liang, G. C. Gardner, and M. J. Manfra, Nature Physics 15, 563 (2019).
[53] M. Henry, S. Oberholzer, C. Strunk, T. Heinzel, K. Esslin, M. Holland, and C. Schönenberger, Science 284, 296 (1999).
[54] I. Neder, N. Ofek, Y. Chung, M. Heiblum, D. Mahalu, and V. Umansky, Nature 448, 333 (2007).
[55] E. Weisz, H. Choi, I. Sivan, M. Heiblum, Y. Gefen, D. Mahalu, and V. Umansky, Science 344, 1363 (2014).