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Abstract
We study junctions between superconductors mediated by the edge states of a quantum-spin-Hall insulator. We show that such junctions exhibit a fractional Josephson effect, in which the current phase relation has a 4π rather than a 2π periodicity. This effect is a consequence of the conservation of fermion parity—the number of electron mod 2—in a superconducting junction and is closely related to the Z_2 topological structure of the quantum-spin-Hall insulator. Inelastic processes, which violate the conservation of fermion parity, lead to telegraph noise in the equilibrium supercurrent. We predict that the low-frequency noise due these processes diverges exponentially with temperature T as T→0. Possible experiments on HgCdTe quantum wells will be discussed.

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Josephson current and noise at a superconductor/quantum-spin-Hall-insulator/superconductor junction

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We study junctions between superconductors mediated by the edge states of a quantum-spin-Hall insulator. We show that such junctions exhibit a fractional Josephson effect, in which the current phase relation has a 4π rather than a 2π periodicity. This effect is a consequence of the conservation of fermion parity—the number of electron mod 2—in a superconducting junction and is closely related to the Z_2 topological structure of the quantum-spin-Hall insulator. Inelastic processes, which violate the conservation of fermion parity, lead to telegraph noise in the equilibrium supercurrent. We predict that the low-frequency noise due these processes diverges exponentially with temperature T as $T \to 0$. Possible experiments on HgCdTe quantum wells will be discussed.

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Proposals for fault tolerant topological quantum computation have motivated intense current interest in finding robust physical systems that host excitations with non-Abelian statistics.[1,2] Recent experiments on the quantum Hall effect have shown encouraging indirect evidence for such excitations,[3,4] but the direct observation of non-Abelions has so far remained elusive. Recently we showed that the proximity effect between a superconductor and a three-dimensional (3D) topological insulator leads to a two-dimensional (2D) interface state that supports non-Abelian Majorana fermions.[5] A first step toward implementing this proposal would be to demonstrate experimentally the topological order responsible for Majorana fermions.

In this Rapid Communication we study Josephson junctions mediated by a 2D topological insulator, known as a quantum-spin-Hall insulator (QSHI).[6–9] We predict such junctions exhibit a *fractional Josephson effect*, which is related to the presence of Majorana fermions. The signature of the fractional Josephson effect is that the current phase relation has a 4π rather than a 2π periodicity. This behavior was first predicted by Kitaev[10] using an idealized model of a one-dimensional (1D) spinless p-wave superconductor. Kwon *et al.*[11] proposed that a related effect can occur at junctions between unconventional 3D superconductors. They argued that it leads to an ac Josephson effect with half the usual Josephson frequency, and that in a weak tunneling limit the Josephson current is carried by electrons rather than Cooper pairs. Michelson *et al.*[12] proposed a related effect in spin active Josephson junctions. The 4π periodicity can occur because the junction has two states with different Josephson currents that are interchanged when the phase is advanced by 2π. At finite temperature inelastic processes can cause transitions between the states, leading to *telegraph noise* in the Josephson current. We will show that in our setup these transitions are forbidden by the local conservation of *fermion parity* (FP), which counts the number of electron mod 2. This leads to an *exponential* suppression of the transition rate at low temperature. This can be probed by measuring the low-frequency current noise $S(\omega \to 0)$, which we predict diverges exponentially at low temperature.

The QSHI is a time-reversal invariant insulating state with a bulk energy gap generated by spin-orbit interactions.[6,8] It has recently been observed in HgCdTe quantum wells.[9] The QSHI is distinguished from an ordinary insulator by a Z_2 topological invariant,[6] which requires the existence of gapless edge states. The edge states form a unique 1D system that is essentially *half* of an ordinary spin degenerate 1D electron gas. In the simplest case it consists of a single band of right moving electrons paired via Kramers theorem with a left moving band with the opposite spin. These states are robust against disorder because time-reversal symmetry prevents elastic backscattering. In the absence of inelastic scattering the edge state transmission is perfect.

Suppose the edge is in intimate contact with an s-wave superconductor. The edge states will become Andreev states, which decay into the superconductor, and may be described with a 1D theory with an induced pairing potential $\Delta=\Delta_0 e^{i\phi}$. Δ_0 depends on the coupling t between the edge and the superconductor.[13] For strong coupling it is of order the bulk gap Δ_{bulk}, while perturbatively it is of order of t^2/Δ_{bulk}. ϕ is the phase of the bulk superconductor. We write[5] $H=\Psi^\dagger \mathcal{H}\Psi/2$, where $\Psi=[(\psi_\uparrow,\psi_\downarrow),(\psi_\downarrow^\dagger,-\psi_\uparrow^\dagger)]$ is expressed in terms of field operators $\psi_{\uparrow(\downarrow)}$ describing the right (left) movers and

$$\mathcal{H} = -iv\tau_z\sigma_z\partial_x - \mu\tau_z + \Delta_0(\cos\phi\tau_x + \sin\phi\tau_y). \tag{1}$$

σ_j are the Pauli matrices acting in the space of right and left movers $\psi_{\uparrow,\downarrow}$ and τ_j are the Pauli matrices which mix the ψ and $\psi^\dagger$ blocks of Ψ. v is the velocity of the edge states, μ is the chemical potential, and we set $\hbar=1$. The eigenstates of Eq. (1) come in pairs at $\pm E$. Due to the redundancy in Ψ, these states are not independent, and the Bogoliubov quasiparticle operators satisfy $\Gamma_{-E}=\Gamma_E^\dagger$.

Equation (1) is similar to Kitaev's model of superconducting spinless electrons in 1D.[10] In Kitaev's model there are zero energy Majorana bound states associated with the *ends* of the sample. In our system, the edge—which is the boundary of the 2D QSHI—cannot have an end. By breaking the time-reversal symmetry, however, a Zeeman field can introduce a mass term into H of the form

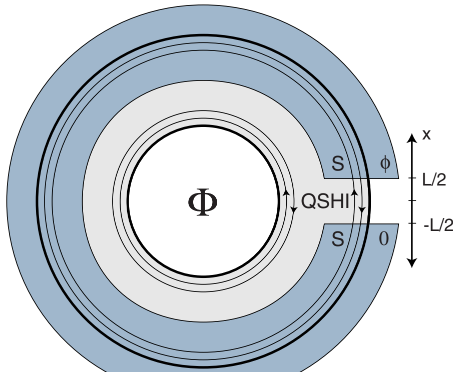

FIG. 1. (Color online) A S/QSHI/S junction in an rf SQUID geometry where the QSHI forms a Corbino disk.

$$V_Z = M\psi^\dagger \sigma_x \psi = M\Psi^\dagger \sigma_x \Psi/2. \tag{2}$$

When $M>\mu$, V_Z opens an insulating gap in the edge state spectrum. V_Z could arise either from an applied magnetic field (as in Ref. 9) or due to proximity to a magnetic material. Zero energy Majorana bound states will exist at the interface between regions with gaps dominated by Δ and M.[5] In the presence of both Δ and M the gap is the smaller of $|\Delta_0 \pm M|$. When $\Delta_0=|M|$ a single band is gapless and for $\Delta_0 \sim |M|$ the low-energy sector of Eq. (1) has the form of a Su-Schrieffer-Heeger model,[14] which has a well known zero energy bound state where $\Delta_0-|M|$ changes sign. The Bogoliubov quasiparticle operator associated with this state is a Majorana fermion, which satisfies $\gamma_0=\gamma_0^\dagger$.

Consider a superconductor/QSHI/superconductor (S/QSHI/S) junction in which the edge states of a QSHI connect two superconductors separated by a distance L. Figure 1 shows an rf superconducting quantum interference device (SQUID) geometry, in which the phase difference across the junction $\phi=(2e/\hbar)\Phi$ is controlled by the magnetic flux Φ. We also assume that the QSHI forms a Corbino disk which circles the flux. This geometry is not essential, but we will see that it has considerable conceptual value. We will also include a Zeeman term in the gap between the superconductors, which will make the connection with Majorana bound states transparent. We emphasize, however, that there will be a nontrivial effect even when this term is absent. To determine the characteristics of the junction we solve the Bogoliubov–de Gennes (BdG) equation $(\mathcal{H}+V_Z)\xi=E\xi$, with

$$\Delta(x) = \Delta_0[\theta(-x-L/2) + e^{i\phi}\theta(x-L/2)],$$

$$M(x) = M_0\theta(x+L/2)\theta(-x+L/2). \tag{3}$$

By enforcing continuity of ψ at $x=\pm L/2$ we determine the spectrum of Andreev bound states in the junction. The calculation is similar to Ref. 11, as well as the theory of superconducting quantum point contacts(SQPCs).[15,16] However, we shall see that there is a fundamental difference with those theories.

Figure 2(a) shows the spectrum as a function of ϕ for $M_0=0$. For $L \lesssim v/\Delta_0$ there is a single pair of bound states $E=\pm\epsilon_0(\phi)$. For $L \ll v/\Delta_0$ our model reduces to the δ function model solved in Ref. 11, where the normal-state transmission probability is $D=1/[1+(M_0 \sinh(\kappa L)/\kappa)^2]$, with $\kappa=\sqrt{M_0^2-\mu^2}$. In that case

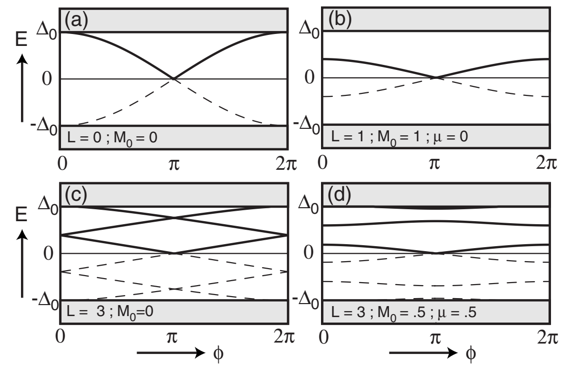

FIG. 2. Spectrum of Andreev bound states in the junction as a function of phase difference ϕ for parameters indicated in each panel. L is in units of v/Δ_0 and M_0 and μ are in units of Δ_0. (a) and (c) are independent of μ.

$$\epsilon_0(\phi) = \sqrt{D}\Delta_0 \cos(\phi/2). \tag{4}$$

Figure 2(b) shows a case where $M_0 \sim \Delta_0$, so the normal-state transmission $D<1$. When $D \ll 1$ there are two weakly coupled Majorana end states at $x=\pm L/2$. When $L>v/\Delta_0$ there will be additional Andreev bound states in the junction with a level spacing of order of v/L. Figure 2(c) shows the case where $L=3v/\Delta_0$ with $M_0=0$, in which time-reversal symmetry requires Kramers degeneracies when $\phi=0$ or π. Figure 2(d) shows the effect of finite M_0 and μ, which lifts most of the degeneracies. However, the crossing at $E=0$ remains and is of special significance.

To understand the crossing consider $E \ll \Delta_0$. The eigenvectors $\xi_{0\pm}$ of Eqs. (1)–(3) with energy $\pm\epsilon_0(\phi)$ define Bogoliubov operators $\Gamma_{0\pm}=\Psi^T\xi_{0\pm}$. Due to particle-hole symmetry, $\Gamma_{0+}=\Gamma_{0-}^\dagger\equiv\Gamma_0$. The low-energy Hamiltonian is thus

$$H = \epsilon_0(\phi)(\Gamma_0^\dagger\Gamma_0 - 1/2) = 2i\epsilon_0(\phi)\gamma_1\gamma_2, \tag{5}$$

where $\gamma_1=(\Gamma_0+\Gamma_0^\dagger)/2$ and $\gamma_2=-i(\Gamma_0-\Gamma_0^\dagger)/2$ are the Majorana operators. For $D \ll 1$ $\gamma_{1,2}$ describe Majorana end states at $x=\pm L/2$ coupled by weak electron tunneling. The crossing at $\phi=\pi$ follows from the destructive interference of the left and right tunneling processes. Equation (5) describes two states distinguished by $N_0\equiv\Gamma_0^\dagger\Gamma_0=0,1$. Mixing these states requires an interaction that changes N_0. Due to the pairing term in Eq. (1), the total charge is not conserved. However, the FP, defined as the number of electron mod 2, is conserved in Eqs. (1)–(3). This forbids the coupling between the two states and protects the crossing at $\epsilon_0(\phi)=0$.

There is a problem, however, with the FP. The junction Hamiltonian [Eqs. (1)–(3)] is invariant under a 2π phase change, but when $\phi\to\phi+2\pi$, the system passes through a single level crossing and can only return to the initial state by a process which changes N_0 by 1. The FP thus apparently changes when $\phi\to\phi+2\pi$. This has to do with the unbounded spectrum as $E\to-\infty$ and reflects a *fermion parity anomaly* similar to the SU(2) anomaly in four-dimensional (4D) field theory.[17] This anomaly is related to non-Abelian statistics. When ϕ advances by 2π, $\gamma_1 \to \gamma_1$ and $\gamma_2 \to -\gamma_2$. In the tunneling limit this can be interpreted as Ivanov's rule[18] for braiding a vortex between the Majorana bound states.

The physical origin of the FP anomaly lies in the topological structure of the QSHI. Consider first the Corbino disk in Fig. 1 without the superconductor. In Ref. 19 we showed that the Z_2 invariant characterizing the QSHI describes the change in the Z_2 "time-reversal polarization" (TRP) when flux $h/2e$ is threaded through the hole. A nonzero TRP specifies a many-body Kramers degeneracy localized at either edge of the disk. Since an odd number of fermions has a Kramers degeneracy, the TRP is precisely the FP. With the superconductor present, start in the ground state at $\Phi=0$. When flux $h/2e$ is threaded through the hole, ϕ advances by 2π *and* a unit of FP is transferred from the inner edge of the disk to the junction on the outer edge. Although Eqs. (1)–(3) is invariant under $\phi \to \phi+2\pi$, the global Hamiltonian, which includes the bulk QSHI, is physically distinct when $\Phi=0$ and $h/2e$.

The local conservation of FP has important consequences for the current and noise in a S/QSHI/S junction. This is most striking near the degeneracy point for $\epsilon_0 \ll \Delta_0$ and $T \ll \Delta_0$. For the remainder of this Rapid Communication we will focus on that regime. We will also consider the limit $L \ll v/\Delta_0$, where there is a single Andreev bound state and Eq. (4) applies although the results can straightforwardly be generalized to the case with multiple Andreev levels provided $T \ll v/L$. In this case, N_0 distinguishes two states, with Josephson currents $I_\pm = \pm I_0$, with

$$I_0(\phi) = \frac{1}{2}\sqrt{D}e\Delta_0 \sin \phi/2. \tag{6}$$

In the absence of transitions that violate local FP conservation there can be no transitions between I_+ and I_-, signaling a fractional Josephson effect.

Elastic-scattering processes can be incorporated into the BdG Hamiltonian from the start and will not lead to violations of the FP. However, at finite temperature, inelastic processes[20,21] can lead to a transition between I_+ and I_- provided an available fermion is present to switch the FP. This could be either due to a thermally excited quasiparticle or due to hopping from a bulk localized state. These processes, however, will be exponentially suppressed at low temperature. On a time scale longer than the switching time the current will thermalize with an average value[11,22]

$$\langle I(\phi)\rangle = I_0(\phi)\tanh \epsilon_0(\phi)/2T. \tag{7}$$

On shorter times, the current will exhibit telegraph noise, as it switches between $I_\pm$.

In order to model the inelastic processes responsible for the telegraph noise we consider the interaction of the Andreev level Γ_0 with a bath of fermions c_n (e.g., quasiparticles) and bosons b_m (e.g., phonons). We thus write

$$H = \epsilon_0\Gamma_0^\dagger\Gamma_0 + \sum_n E_n c_n^\dagger c_n + \sum_m \omega_m b_m^\dagger b_m + \sum_{mn}[(V^1_{nm}c_n^\dagger b_m + V^2_{nm}c_n b_m^\dagger)\Gamma_0 + \text{H.c.}]. \tag{8}$$

Here $E_n, \omega_n > 0$, and we have ignored terms which create (or annihilate) both fermions and bosons. The transition rates $\tau_\pm^{-1}(\epsilon_0, T)$ between the states N_0 and $N_0 \pm 1$ follow from Fermi's golden rule. For $\epsilon_0, T \ll \Delta_0$ we find

$$\tau_\pm^{-1} = e^{\mp\epsilon_0/2T}[w_1(T)e^{\epsilon_0/2T} + w_2(T)e^{-\epsilon_0/2T}], \tag{9}$$

where

$$w_{1,2}(T) = 2\pi\sum_{n,m} e^{-E_n/T}|V^{1,2}_{nm}|^2\delta(E_n - \omega_m). \tag{10}$$

If either the Zeeman term vanishes ($M_0=0$) or the system is symmetric under $x \to -x$, then $w_1(T)=w_2(T)\equiv w(T)$. We will assume this below although the results are only slightly modified otherwise. $w(T)$ depends on the dominant source of fermions, which we take to be either thermally activated quasiparticles or Mott variable range hopping from bulk localized states,

$$w(T) \propto \begin{cases} e^{-\Delta_0/T} & \text{quasiparticles} \\ e^{-(T_0/T)^{1/3}} & \text{hopping}. \end{cases} \tag{11}$$

T_0 depends on the density of states and localization length, and we assume the hopping is 2D.

The transition rate is exponentially suppressed for $T \to 0$. At sufficiently low temperature the resulting telegraph noise could be observed in the time domain. At higher temperature there is a signature in the noise spectrum $S(\omega)$. We determine $S(\omega)$ semiclassically by solving a kinetic equation for the probability $p(t)$ that $N_0=1$.[20,21] This has the form $dp/dt = -(p-\bar{p})/\tau$, where $\tau^{-1}=\tau_+^{-1}+\tau_-^{-1}=4w\cosh^2 \epsilon_0/2T$. $\bar{p}=(1+\exp \epsilon_0/T)^{-1}$ follows from the detailed balance condition $\tau_+/\tau_-=e^{\epsilon_0/T}$. Temporal correlations in $I(t)$ decay exponentially on a time scale w^{-1}, and the noise spectrum $S(\omega) = 2\int_{-\infty}^{\infty} e^{i\omega t}\langle I(t)I(0)\rangle$ is given by[20,23]

$$S(\omega) = \frac{4I_0^2}{\cosh^2 \epsilon_0(\phi)/2T}\frac{\tau}{1+\omega^2\tau^2}. \tag{12}$$

In the zero-frequency limit we have

$$S(\omega \to 0) = \frac{I_0^2}{w(T)\cosh^4 \epsilon_0(\phi)/2T}. \tag{13}$$

For $D=1$, these results are similar to the theory of a SQPC.[20–23] However, the current in Eq. (6) is half the value of a perfect single-channel SQPC. A SQPC is similar to two copies of a S/QSHI/S junction. This leads to a fundamental difference because in the SQPC there is no conservation law to prevent scattering between the $\pm I_0$ states, which can occur via low-energy processes that transfer an electron between the two pairs. Elastic backscattering in the SQPC leads to an avoided crossing of the states near $E=0$, so the Andreev states carry no current at $\phi=\pi$. It is also of interest to compare with the theory of Ref. 11. In that work, multichannel junctions were considered. Independence of the different channels requires translational symmetry, so impurity scattering will lead to the violation of the conservation FP within a given channel. In Ref. 12 it was argued that a 4π periodicity of the Josephson current is possible for a spin active junction with mirror symmetry although similar to the SQPC perfect symmetry is required. In addition, FP protected $E=0$ level crossings can occur in that system, which could exhibit telegraph noise. However they occur at two distinct phases ϕ and $2\pi-\phi$ and are not topologically guaranteed. The low-temperature behavior predicted by Eqs. (11) and (13) is unique to the S/QSHI/S junction and is a signature of the FP anomaly.

We now briefly consider junctions at finite voltage bias. There are two cases, depending on M_0. For $M_0=0$, the perfect edge state transmission causes the Andreev levels to merge with the continuum. This leads to a finite dc current, which for $eV \ll \Delta_0$ can be understood semiclassically in terms of multiple Andreev reflections.[23] For $w(T) \ll eV \ll \Delta_0$, the current is $I(V)=(2/\pi)I_c \operatorname{sgn} V$, where $I_c=\sqrt{D}e\Delta_0/2$. For $M_0>0$, there is an energy gap δ separating the Andreev levels from the continuum, as in Figs. 2(b) and 2(d). For $w(T) \ll eV \ll \delta$ there will be a fractional ac Josephson current with frequency $eV/\hbar$.[11] For $eV \sim \delta$ Landau-Zener tunneling processes through δ will lead to a damping of the ac Josephson current as well as a finite dc current.

We close by discussing the feasibility of experiments using the QSHI in HgCdTe quantum wells,[8,9] which has a bulk gap of order of 20 meV.[24] The desired geometry would be similar to Ref. 25, where a 2D InAs quantum well was contacted with Nb. The gap Δ_0 will depend on the contact, and if optimized could be of order the bulk gap of the superconductor. The electrode that is currently used in HgCdTe quantum wells is indium, which becomes an s-wave superconductor below $T_c=3.4$ K.[26] Using $v=3.6$ eV Å (Ref. 24) and $\Delta_0=0.1$ meV we find $L \lesssim v/\Delta_0 \sim 3$ μm sets the scale for having a single Andreev level. The simplest experiment would be to study a single-current-biased junction, which is predicted to have a critical current $I_c=e\Delta_0/2 \sim 10$ nA. Measuring the equilibrium telegraph noise at $\phi \sim \pi$ requires an inductive measurement on a rf SQUID.[16] The physics at $M_0 \neq 0$ requires a magnetic field in the junction region. An appropriately aligned field induces a gap $B \times (3.1\ \text{meV/T})$ (Ref. 24) in the edge states, so a field of order of $0.03T$ could suppress the normal-state transmission D as well as the magnitude of the Josephson current.

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