Robust Transport Signatures of Topological Superconductivity in Topological Insulator Nanowires

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Finding a clear signature of topological superconductivity in transport experiments remains an outstanding challenge. In this work, we propose exploiting the unique properties of three-dimensional topological insulator nanowires to generate a normal-superconductor junction in the single-mode regime where an exactly quantized $2e^2/h$ zero-bias conductance can be observed over a wide range of realistic system parameters. This is achieved by inducing superconductivity in half of the wire, which can be tuned at will from trivial to topological with a parallel magnetic field, while a perpendicular field is used to gap out the normal part, except for two spatially separated chiral channels. The combination of chiral mode transport and perfect Andreev reflection makes the measurement robust to moderate disorder, and the quantization of conductance survives to much higher temperatures than in tunnel junction experiments. Our proposal may be understood as a variant of a Majorana interferometer which is easily realizable in experiments.

A topological superconductor is a proposed novel phase of matter with exotic properties like protected boundary states and emergent quasiparticles with non-Abelian statistics. If realized, these superconductors are expected to constitute the main building block of topological quantum computers [1]. The prototypical example of this phase, the $p$-wave superconductor, has proven to be difficult to find in nature, with superconducting Sr$_2$RuO$_4$ and, indirectly, the $\nu=5/2$ fractional quantum Hall state among the very few conjectured candidates. While many experiments have been suggested and performed on these systems, evidence for their topological properties remains elusive. However, the recent realization that a $p$-wave superconductor need not be intrinsic, but can alternatively be engineered with regular $s$-wave superconducting proximity effect in strongly spin-orbit coupled materials [2–4], has opened a promising new path in the search for topological superconductivity.

A class of these new topological superconductors is predicted to be realized in one-dimensional (1D) systems with broken time-reversal symmetry [5]. These systems feature Majorana zero-energy end states, which are responsible for a fundamental transport effect known as perfect (or resonant) Andreev reflection [6]. When a single propagating mode impinges on a 1D superconductor, unitarity and particle-hole symmetry require that it must be either reflected as an electron, or Andreev-reflected as a hole. In the first case, the normal-superconductor (NS) conductance vanishes, but in the second a Cooper pair is transferred across the junction resulting in an exactly quantized zero-bias conductance of $2e^2/h$. The second case occurs when the superconductor is topological and can be intuitively understood as resonant transport mediated by the Majorana end states [6, 8]. This difference in conductance in the single-mode regime can be taken as the defining feature of a 1D topological superconductor [7, 8], thereby potentially allowing for its direct identification.

A prominent example of a 1D topological superconductor is realized in semiconducting quantum wires in the presence of a magnetic field [11, 12]. Recent transport experiments with InSb and InAs wires aimed to demonstrate the existence of this phase have reported a finite zero-bias conductance across a NS junction [13, 14], but the predicted quantization has so far remained a challenge to observe. A possible reason is that these wires typically host several modes [15–19] and fine tuning the chemical potential to the single mode regime can be difficult. In the presence of several modes, either a tunnel...
dimensional topological insulators (TI) such as Bi$_2$Se$_3$ starting from an alternative route to 1D topological insulator. This effect should therefore have a robust, easy to tune in and out of the topological phase.

In this work, we propose to realize such a junction starting from an alternative route to 1D topological superconductivity, recently proposed by Cook and Franz [20], based on the use of nanowires made from three dimensional topological insulators (TI) such as Bi$_2$Se$_3$. In the nanowire geometry, the 2D surface states of a TI are resolved into a discrete set of modes, with the special property that when a parallel flux of $h/2e$ threads the wire, the number of modes within the bulk gap is always odd [21][23]. When a superconducting gap is induced on the surface via the proximity effect, this guarantees that the system becomes a topological superconductor [5][24]. An NS junction can then be built by proximitizing only the perimeter with vorticity $n_v$, while the superconducting part can be tuned in and out of the topological phase at will with the in-plane flux $\eta$.

In addition, our design simultaneously allows to drive the normal region of the junction into the single-mode regime by exploiting the unique way TI surface states respond to magnetic fields [25][29]. When a perpendicular field is applied to the normal part of the wire, its top and bottom regions become gapped, while counter-propagating chiral edge states form in the side regions in between. These are analogous to quantum Hall edge states and are protected from backscattering due to their spatial separation. The resulting NS junction, shown in Fig. 1, has a single chiral mode reflecting from the superconductor, and is ideal for probing conductance quantization. Moreover, all of its components are readily available, as both surface transport in TI nanowires [30][33] and the contacting of bulk TI with superconductors [34][36] have already been demonstrated experimentally. In the remainder of this paper, we provide a detailed study of the transport properties of this system, demonstrating that conductance quantization is indeed achievable under realistic conditions, and discuss the advantages of our setup over other proposals.

To model the proposed device, we consider a rectangular TI nanowire of height $h$ and width $w$ (cross section area $A = hw$ and perimeter $P = 2h + 2w$). The surface of the wire is parametrized with two coordinates $(x, s)$, where $s$ is periodic $s \in [0, 2\pi]$ and goes around the perimeter of the wire, while $x$ goes along its length. We first consider a magnetic field parallel to the wire, $\vec{B} = (B_0, 0, 0), \text{described with the gauge choice } \vec{A} = B_0(0, -z/2, y/2)$. The dimensionless flux through the wire is $\eta = B_0 A/(h/e)$. In this simple case the effective theory for the surface states is the same as for a cylindrical wire [21][22], with $s$ playing the role of the azimuthal angle

$$H_0 = -i\psi_F [\sigma_x \partial_x + \sigma_y (2\pi/P)(\partial_s + i\eta)],$$

where we set $\hbar = 1$ and take $v_F = 330$ meV nm [37]. The Zeeman coupling to the magnetic field produces no qualitative differences in the physics we describe [35]. The wavefunctions satisfy antiperiodic boundary conditions in $s$ due to the curvature-induced $\pi$ Berry phase [21][22]. The eigenfunctions of $H_0$ thus have the form

$$\psi_{k,n}(x, s) = e^{i k x} e^{i\eta s} \chi_{k,n},$$

with half-integer angular momentum $l_n = n - 1/2$ where $n \in \mathbb{Z}$. The spectrum is

$$E_{k,n} = v_F \sqrt{k^2 + (2\pi/P)^2(n + \eta)^2},$$

and is depicted in Fig. 2. For $\eta = 0$ all modes are doubly degenerate, while for $\eta = 1/2$ the number of modes is always odd because the $n = 0$ one is not degenerate.

By bringing the wire into contact with an s-wave superconductor [20], as shown in Fig. 1 an s-wave pairing potential $\Delta$ is induced due to the proximity effect. The Bogoliubov-de Gennes Hamiltonian can be written as $H = \frac{1}{2} \Psi^\dagger \mathcal{H} \Psi$ with

$$\mathcal{H} = \begin{pmatrix} H_0 & \Delta(s) \\ \Delta^*(s) & -T^{-1}H_0 T \end{pmatrix},$$

where $\Psi = (\psi, \psi^\dagger, -\psi^\dagger)$ is a Nambu spinor. The induced pairing potential takes the form $\Delta(s) = \Delta_0 e^{-i n_v s}$, where we have allowed the phase of $\Delta$ to wind around the perimeter with vorticity $n_v$. For $\eta = 0$ the ground state has $n_v = 0$. Around $\eta = 1/2$, however, it should be energetically favorable for $\Delta$ to develop a vortex. In an actual experiment, $n_v$ is then expected to jump abruptly as $\eta$ is ramped continuously from zero to $1/2$ [39]. For $\eta$ around $1/2$ and in the presence of a vortex, the nanowire becomes a topological superconductor for any $\mu$ within the bulk gap [20].

The presence of the vortex is essential in order to observe perfect Andreev reflection in our setup [40]. To see this, consider the Hamiltonian in Eq. 4 in the presence of a NS interface at $x = 0$, with $n_v$ vortices in the superconducting part. Introducing Pauli matrices $\tau_i$ acting in Nambu space

$$\mathcal{H}^{(n_v)} = \begin{pmatrix} -i\sigma_x \partial_x + \sigma_y (2\pi/P)(\partial_s + \eta) & 0 \\ 0 & -i\sigma_x \partial_x + \sigma_y (2\pi/P)(\partial_s + \eta) - \Delta_0 e^{-i \eta s} \tau_z \end{pmatrix}$$

For $n_v = 0$, electron states in the normal part have finite angular momentum $l_n$, see Eq. 2, while hole states have angular momentum $-l_n$, independently of the value of $\eta$. Since angular momentum must be conserved upon
reflection, a single incoming electron can never be reflected as a hole. For $n_v = 1$ rotational invariance appears to be broken by the pairing term, but is explicitly recovered after the gauge transformation $\Psi \rightarrow e^{i\tau_3/2}\Psi$, which shifts $\eta \rightarrow \eta - 1/2$. This transformation also changes the boundary conditions to periodic, such that angular momenta take integer values $l'_n = n$. As a result, the $n = 0$ electron state now has the same angular momentum as its conjugate hole state and can be reflected into it. The vortex provides the extra unit of angular momentum required to reflect one into the other.

The NS conductance of the junction is computed from the Andreev reflection matrix, evaluated separately for electron states, labelled now by $e$, and for hole states, labelled by $h$. Normalization is chosen such that all propagating states carry the same current, $J_x = \langle \psi | \sigma_x | \psi \rangle = 1$. These are matched to the evanescent states in the superconductor $\psi^{S+}_n$ and $\psi^{S-}_n$ by imposing continuity of the wavefunction at the junction (dropping the label $n$ for ease of notation)

$$\psi^{e-} + r_{ee}\psi^{e+} + r_{eh}\psi^{h-} = a\psi^{S+} + b\psi^{S-},$$

(6)

$$\psi^{h-} + r_{hh}\psi^{h+} + r_{eh}\psi^{e-} = a'\psi^{S+} + b'\psi^{S-}.$$  

(7)

The reflection matrix is defined as $r = (r_{ee} \, r_{eh})$, and is both unitary and particle-hole symmetric. The conductance is given by $G_{NS} = 2e^2/h \int r_{eh} d\epsilon$, where the trace sums over all propagating modes. The resulting $G_{NS}$ for $\Delta_0 = 0.25$ meV are shown in Fig. 2(b). When $\eta = 1/2$, $n_v = 1$ and in the range $\mu < \pi/P$, a single mode is reflected from a topological superconductor resulting in a conductance of $2e^2/h$.

The conditions to observe conductance quantization in this setup are not optimal yet, mainly because the chemical potential has to be tuned into a small gap $\pi/P$. This limitation can be overcome to a large extent by the addition of a perpendicular field. Consider the Hamiltonian of the normal wire with $\vec{B} = (B_\parallel, B_\perp, 0)$ and a vector potential $\vec{A} = B_\perp (z, 0, 0) + B_\parallel (0, -z/2, y/2)$ such that translational invariance is still preserved in the $x$ direction

$$H = \sigma_x [-i\partial_x + eA_x(s)] + \frac{2\pi}{P} \sigma_y (-i\partial_s + \eta),$$

(8)

where the vector potential in the surface coordinates is given by

$$A_x(s) = B_\perp P \begin{cases}
\frac{s}{2\pi} & \frac{1}{4} < \frac{s}{2\pi} < \frac{3}{4} \\
\frac{s}{2\pi} + \frac{3}{4} & \frac{3}{4} < \frac{s}{2\pi} < 1 \\
\frac{s}{2\pi} & \frac{1}{4} \leq \frac{s}{2\pi} < \frac{3}{4} \\
\frac{s}{2\pi} + \frac{3}{4} & \frac{3}{4} \leq \frac{s}{2\pi} < 1 \end{cases},$$

(9)

with $r = \frac{w}{w + h}$. The profiles of $A_x(s)$ and $B_\perp$ along the $s$ direction are shown in the inset of Fig. 2(b). Since rotational symmetry is broken, the different $n$ modes are mixed. In the angular momentum basis, Eq. (2), the Hamiltonian is $H = \sum_{n, n'} (-\Delta_0) \chi_{k, n'} \chi_{k, n}$. For $\eta = 0$, $n_v = 0$, $eV = 0$ as a function of $\mu$, $\mu = 0$, and $\eta = 1/2$, $n_v = 1$ (full line) c) The same for $B_\parallel = 2T$.

FIG. 2. a) The spectrum of a wire of dimensions $h = 40$ nm and $w = 160$ nm for $B_\perp = 0$ (center) and $B_\perp = 2T$ (right). Note that in the last case the spectrum is independent of $\eta$. b) In the range $\mu < \pi/P$, a single mode is reflected from a topological superconductor resulting in a conductance of $2e^2/h$.
α' = 1, ..., N_{ev}, with N_{ev} + N_{prop} = 2N. Both propagating and evanescent momenta and wavefunctions are obtained efficiently from the transfer matrix of the normal part Hamiltonian [58]. We assume that B_⊥ is completely screened in the superconducting part of the wire (see Fig. 1), so that the eigenstates in this region remain the same as before. Continuity of the wavefunctions at the interface implies the matching condition

\[ \psi^{-}_\alpha + \sum_{\beta=1}^{N_{prop}} \left[ (r_{ee})_{\alpha\beta} \psi^+_{\beta} + (r_{he})_{\alpha\beta} \psi^0_{\beta} \right] = \sum_{n=1}^{N_{ev}} \left[ c_{\alpha n} \psi^+_{n} + d_{\alpha n} \psi^{-}_{n} \right] \]

For every value of \( \alpha \), we project into angular momentum states with \( n = -N, ..., N \), and since the spinors have four components (spin and particle-hole degrees of freedom) this yields a system of 8N equations with 2N_{prop} + 2N_{ev} + 4N = 8N coefficients. The system is solved numerically, and the conductance obtained is shown in Fig. 2(c). In the single-mode regime, at zero flux and \( n_\nu = 0 \) we have \( G_{\text{NS}} = 0 \), but at \( \eta = 1/2 \) and \( n_\nu = 1 \) (when the superconductor is topological), we have \( G_{\text{NS}} = 2e^2/h \) as expected.

The quantization of \( G_{\text{NS}} \) can be understood intuitively in terms of a 1D low-energy model, depicted in Fig. 1(b), similar to the one describing the Majorana interferometer proposed in Refs. [9] and [10]. In this model, an incoming chiral mode leaving the source is split into two Majorana modes that appear at the interface between the superconductor and the regions with finite \( B_\perp \). In the absence of a vortex in the superconductor the two Majoranas recombine as an electron on the other side of the wire and come back to the source through the channel of opposite chirality, yielding \( G_{\text{NS}} = 0 \). However, if a vortex is present, the two Majoranas accumulate a relative phase of \( \pi \) and recombine as a hole, while a Cooper pair is transferred to the superconductor, yielding \( G_{\text{NS}} = 2e^2/h \).

The quantization of the conductance in our setup is expected to be robust to disorder to some extent, because transport in the normal part is mediated by spatially separated chiral modes. In order to test this robustness explicitly, we introduce disorder into the Hamiltonian of a normal wire in the presence of \( B_\perp \), and compute the two terminal conductance \( G_N \) of a finite size wire numerically, following the method of Ref. [22]. The disorder potential has a correlator \( \langle V(r)V(r') \rangle = g^2/2\pi \xi_0^2 e^{-|r-r'|^2/2\xi_0^2} \), with \( \xi_0 \) the disorder correlation length and \( g \) a dimensionless measure of the disorder strength. Our data is obtained by averaging over \( 10^3 \) disorder configurations. The results are shown in Fig. 3. We observe that in the single-mode regime, the conductance of the normal wire indeed remains quantized to \( e^2/h \) in the presence of moderate disorder, as long as the chemical potential is not very close to zero. A full characterization of the effects of disorder will be presented in a future work [38].

**Discussion** - In our setup, a quantized conductance can be obtained with both \( B_\perp = 0 \) and finite \( B_\perp \), but the latter case has several advantages that are worth stressing. First, the single-mode regime remains accessible for chemical potential values ranging up to values of the order of the cyclotron frequency \( \omega_c \), rather than the finite size gap \( \pi/P \), as Fig. 2(c) demonstrates. Second, chiral mode transport in the normal part is robust against finite disorder due to spatial separation of counter propagating chiral modes. Third, the spectrum of the normal part in the presence of \( B_\perp \) becomes independent of \( B_\parallel \), which in turn affects only the superconducting part. \( B_\parallel \) thus becomes an independent knob driving the transition from a trivial to a topological superconductor, while the chiral modes remain intact. In this case, measuring \( G_{\text{NS}} = 0 \) would represent a genuine consequence of reflection from a trivial superconductor, as opposed to the \( B_\perp = 0 \) case where this value of \( G_{\text{NS}} \) could result simply from an insulating normal part, see Fig. 2(b).

Our proposal realizes a version of the Majorana interferometer with some important differences. In our setup, instead of contacting the two chiral modes separately, see Fig. 1(c), the source electrode contacts both channels and the superconductor is the drain [9], see Fig. 1(d). In addition, the original proposals use ferromagnets and a finite superconducting island to create the Majorana modes, while our setup uses a bulk superconductor and a homogeneous magnetic field [11], making it experimentally more feasible. Despite these differences, the finite voltage and finite temperature behavior of \( G_{\text{NS}} \) in our system will be similar to those in Refs. [9] and [10]. This introduces an important advantage to our setup over current semiconducting wires, where the temperatures required to observe quantization of the conductance are of the order of mK. In our setup, the limiting temperature
is determined by the proximity induced gap \[10\]. Assuming \(\Delta_0 \approx 0.1 - 0.25 \text{ meV} \[13\] \[14\] \[15\] this corresponds to 1-3 K, an improvement of several orders of magnitude.

Finally, we note that screening \(B_1\) in the SC region requires the use of a superconductor with a high critical field so that NS transport may be measured in the regime where chiral states form. For example, the superconductor could be a Ti/Nb/Ti trilayer as the one used to contact an InSb nanowire in the experiment in Ref.\[43\] which was estimated to have \(H_{c1} = 2.5\text{T}\).

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