Nonlocal Hanbury–Brown–Twiss interferometry and entanglement generation from Majorana bound states

Sougato Bose\textsuperscript{1,4} and Pasquale Sodano\textsuperscript{2,3}

\textsuperscript{1} Department of Physics and Astronomy, University College London, Gower Street, London WC1E 6BT, UK
\textsuperscript{2} Department of Physics and Sezione INFN, University of Perugia, Via A Pascoli, 06123 Perugia, Italy
\textsuperscript{3} Perimeter Institute of Theoretical Physics, Waterloo, ON, N2L 2Y5, Canada
E-mail: s.bose@ucl.ac.uk

New Journal of Physics 13 (2011) 085002 (9pp)
Received 6 March 2011
Published 3 August 2011
Online at http://www.njp.org/
doi:10.1088/1367-2630/13/8/085002

Abstract. We show that a one-dimensional device supporting a pair of Majorana bound states (MBS) at its ends can produce remarkable Hanbury–Brown–Twiss-like interference effects between well-separated Dirac fermions of pertinent energies. We find that the simultaneous scattering of two incoming electrons or two incoming holes from the MBS leads exclusively to an electron–hole final state. This ‘anti-bunching’ in electron–hole internal pseudospin space can be detected through current–current correlations. Further, we show that, by scattering appropriate spin-polarized electrons from the MBS, one can engineer a non-local entangler of electronic spins for quantum information applications. Both the above phenomena should be observable in diverse physical systems enabling us to detect the presence of low-energy Majorana modes.

\textsuperscript{4} Author to whom any correspondence should be addressed.
1. Introduction

Quantum indistinguishability has striking manifestations when two identical particles are brought together at a beam splitter. For example, two bosons in identical states would ‘bunch’ together when exiting a beam splitter purely due to interference effects [1]. Two fermions, on the other hand, would exit separately or ‘anti-bunch’ [2]. These effects are indeed an instance of the celebrated Hanbury–Brown–Twiss effect, which has also been tested recently with helium atoms [3]. The same quantum indistinguishability is used for the production of entangled photons [4] and can also be used to entangle generic massive particles [5]. Of course, all these effects can occur only when the particles are brought together spatially, for instance, at a beam splitter. It is thereby interesting to look for settings where rather well-separated identical particles could manifest such phenomena.

In this paper we report on the possibility of engineering a non-local beam splitter enabling the above class of phenomena for distant charged fermions. Here, by ‘non-local’ we mean spatially extended. Going beyond the usual two-particle interference in orbital/momentum space, here one finds a Hanbury–Brown–Twiss effect in the electron–hole internal pseudospin space. This is enabled by Majorana mid-gap low-energy modes that transform between electrons and holes [6], effectively making them indistinguishable in a scattering experiment. This Hanbury–Brown–Twiss effect is thereby a detector of the Majorana modes.

Recently, low-energy Majorana (neutral charge self-conjugated fermion) modes located at the edges of linear devices were predicted to induce non-local phenomena [7–9]. Indeed, there are a variety of platforms to realize such devices: for example, a quantum wire immersed in a p-wave superconductor [7, 10], cold-atomic systems mimicking p-wave superconductors [11], topological insulator–superconductor–magnet structures [8, 12, 13] and potentially also semiconductor systems [14, 15]. Evidence for their non-local nature includes distance-independent tunneling [7], crossed Andreev reflection [8] and teleportation-like coherent transfer of a fermion [9]. Finally, they may be easily manipulated [6] and are relevant excitations also in conventional superconductors [16]. To date, the primary application envisaged for these fermions has been topological quantum computation [17]. As the second key result of this paper, we will show another use of these modes, namely that Majorana bound states...
Figure 1. The non-local beam splitter and the electron spin entangler. The MBS are shown as empty ellipses 1 and 2. One specific realization where MBS occur at the boundaries between magnets (M) and superconductors (SC) deposited on quantum spin Hall insulators is depicted, although our results hold more generally. Incoming and outgoing particles are indicated by arrows and may, in practice, be tunneled in/out by STM tips or electron pumps acting as leads.

(MBS) could be used to engineer entanglement between the spins of well-separated particles, a pivotal resource in quantum information.

This paper is organized as follows. In section 2, we consider the scattering of Dirac fermions off the edge MBS of energy $E_M$ in the spinless model investigated in [8]. Here, we show that when the energy of the incoming fermions is nearly resonant with $E_M$, the edge MBS induce a beam splitting process that acts as an equally weighted four-port beam splitter, with ports corresponding to both spatial and electron–hole isospin states. In section 3, we show that with two incident Dirac fermions, this allows for fermion antibunching in the pseudospin space, which has holes and electrons as its two states. This is one of the central results of this paper. In section 4, we determine the signature of this fermion antibunching in the zero-frequency spectral density of the current fluctuations in the leads. Section 5 generalizes the results of section 2 by accounting for the spins of the fermions. In section 6, we show that the edge MBS allow for generating entanglement between the spins of distant electrons only by pertinently choosing the polarizations of the incoming fermions. In section 7, we analyze a few condensed matter settings where our findings may be helpful in detecting the presence of MBS; section 8 is devoted to a few remarks on our results. For a reader interested more in the logical steps leading to our central results, rather than the full technical details, we recommend focusing primarily on sections 3 and 6, taking the relevant scattering matrices from the sections immediately preceding them.

2. Nearly resonant electron (hole) scattering from Majorana edge states

We consider a one-dimensional device supporting two weakly coupled MBS at its ends as shown in figure 1. The MBS are labeled as 1 and 2 and schematically shown as empty ellipses in the figure. As the separation between the MBS increases, their energy $E_M$ decreases exponentially [7]. For the sake of clarity, we will first show how this device produces Hanbury–Brown–Twiss-like interference effects between spatially separated Dirac fermions in the spinless models investigated in [7, 10]; later, we show how all the results are valid for more realistic spinfull physical settings [12, 14, 15]. The Hamiltonian describing the weak coupling between the MBS $\gamma_1$ and $\gamma_2$ is given by

$$H_M = iE_M\gamma_1\gamma_2,$$

(1)
where \( \gamma_j \) are Majorana operators defined by \( \gamma_j = \gamma_j^\dagger \) and satisfying \( \gamma_j \gamma_k + \gamma_k \gamma_j = 2 \delta_{kj} \) (in our definition, the \( \gamma_j = c_j + c_j^\dagger \) in terms of Dirac fermionic operators \( c_j \)).

Leads, also labeled as 1 and 2, are connected to the device as shown in figure 1, allowing for the scattering of Dirac fermions (electrons or holes) from each of the MBS; we further assume that lead 1 is coupled only to bound state 1 and lead 2 is coupled only to bound state 2. The Hamiltonian describing the leads need not be specified at this stage since—for the evaluation of the \( S \) matrix—the leads may be effectively removed by introducing complex embedding potentials [18]. Using the approach developed in [18, 19], one may then write the unitary scattering matrix \( \mathcal{S} \) in a form that is formally independent of the model used to describe the leads; for the scattering of fermions from the MBS located at the ends of the quantum wire shown in figure 1, the \( \mathcal{S} \) matrix has been computed in [8] as

\[
\mathcal{S}(E) = 1 + 2\pi i W^\dagger (H_M - E - i\pi WW^\dagger)^{-1} W,
\]

where \( W \) is a rectangular matrix

\[
\begin{pmatrix}
  w_{11} & 0 & w_{12}^* & 0 \\
  0 & w_{21} & 0 & w_{22}^*
\end{pmatrix}
\]

in the basis \( \{|e_1\}, |e_2\rangle, |h_1\rangle, |h_2\rangle \), with \( |e_j\rangle \) and \( |h_j\rangle \) representing an electron and a hole in the lead \( j \). \( W \) describes the coupling of the scatterer \( (H_M) \) to the leads and \( E \) is the energy of the incident electrons/holes. The entries \( w_j \) of the \( W \) matrix are related to the couplings to the leads \( \Gamma_j = 2\pi w_j^2 \) [8].

For our purposes, it is convenient to assume that \( E \gg \Gamma_j \), as well as \( E \approx E_M \) (i.e. the energies of the incoming Dirac fermions are tuned to be nearly resonant with the Majorana coupling energy). Under these circumstances, the \( \mathcal{S} \) matrix simplifies to

\[
\mathcal{S} = \frac{1}{2} \begin{pmatrix}
  1 & -i & 1 & -i \\
  i & 1 & i & -1 \\
  -1 & -i & 1 & -i \\
  i & -1 & i & 1
\end{pmatrix}
\]

where the basis is, again, \( \{|e_1\}, |e_2\rangle, |h_1\rangle, |h_2\rangle \) ). Note that this regime is different from the one considered by Akhmerov and coworkers [8], where only the terms corresponding to crossed Andreev reflection (i.e. \( \langle h_2 | S | e_1 \rangle \) and \( \langle h_1 | S | e_2 \rangle \)) are maximized. Here, we work in a regime where all the entries of \( \mathcal{S} \) have the same magnitude. It is the implications of this scattering matrix \( \mathcal{S} \) of equation (3) that we work out in this paper. It is this \( \mathcal{S} \) that enables both the non-local Hanbury–Brown–Twiss interferometry in isospin space and the non-local entanglement generation.

Let us first illustrate the action of the above \( \mathcal{S} \) matrix by describing what happens to electrons tunneling in from one end of the one-dimensional device. If, at time \( t \), a single electron tunnels into the Majorana mode located at site 1, i.e. the incoming state is \( c_1^\dagger |0\rangle \), it transforms, under \( \mathcal{S} \), to

\[
c_1^\dagger |0\rangle \xrightarrow{\text{MBS}} \frac{1}{2} (c_1^\dagger + ic_2^\dagger - d_1^\dagger + id_2^\dagger) |0\rangle,
\]

where \( c_j^\dagger (d_j^\dagger) \) creates an electron (hole) at site \( j \). In equation (4), we have used MBS above the arrow to indicate that MBS are responsible for the process. Since the transformation (4) is equivalent to a four-port beam splitter, with MBS inducing the beam splitting process, one
can equally well take MBS to stand for the ‘Majorana beam splitter’. Equation (4) implies that an incoming electron has one-fourth probability of coming out of each site as an electron or a hole. If another electron scatters at a different time \( t' \) on the Majorana mode located at position 2, it will also scatter with exactly the same probabilities for the four possible outcomes. The joint probability for two incoming electrons to exit as two electrons or two holes (whichever the output port) would thus be \( \frac{1}{4} \).

3. The Hanbury–Brown effect in pseudospin space

We will now show that when \( t = t' \), i.e. simultaneous scattering, two-particle interference can take place so that the probability of two electrons or two holes exiting is completely suppressed. By \( t = t' \) we mean that the wavepackets of the two incoming electrons (holes) are large enough so that their time of arrival cannot be distinguished when one observes them after the scattering.

When two electrons scatter simultaneously, one at site 1 and the other at site 2, one has

\[
c_1^\dagger c_2^\dagger |0\rangle \xrightarrow{\text{MBS}} \frac{1}{2}(c_1^\dagger + ic_2^\dagger - d_1^\dagger + id_2^\dagger) \cdot \frac{1}{2}(-ic_1^\dagger + c_2^\dagger - id_1^\dagger - d_2^\dagger)|0\rangle
\]

\[
= \frac{1}{4}(ic_1^\dagger d_1^\dagger - c_1^\dagger d_2^\dagger + ic_2^\dagger d_1^\dagger + ic_2^\dagger d_2^\dagger)|0\rangle.
\]

In the last step of equation (5), we have used \( d_j^\dagger (E) = c_j (-E) \) (which effectively embodies the indistinguishability of an electron and a hole), where \( E \) is the energy. From equation (5), one sees that the probability of two outgoing electrons (holes) after the scattering is zero. Exactly the same holds when two holes scatter simultaneously at leads 1 and 2. This is an interference effect in the same sense as the anti-bunching of fermions at a normal two-port beam splitter, where fermions cannot exit through the same port. Instead of being in the spatial channels, here the anti-bunching is in the internal pseudospin space which has particle and hole as its two states. The unitary conversion of an electron to a hole is, per se, not surprising in view of [6].

Of course, in a practical realization, the condition \( E \sim E_M \) required for obtaining the scattering matrix \( S \) of equation (3) may not be exactly met. To see the effect of an energy mismatch, we denote by \( \delta E \) the amount by which \( E \) deviates (either positively or negatively) from \( E_M \); this deviation is, however, assumed to be much lower than \( E_M \) itself (i.e. \( \delta E \ll E_M \)). Without assuming \( \delta E \ll E_M \), one may end up in qualitatively different regimes: for example, for \( \delta E \) comparable to \(-E_M \), one reaches the regime of [8] of only crossed transmission. For \( \delta E \ll E_M \), the scattering matrix as a function of \( \delta E \) is given by

\[
S_{\delta E} = -\frac{i\Gamma}{\delta E + i\Gamma} S + \frac{\delta E}{\delta E + i\Gamma} I,
\]

where \( \Gamma = \Gamma_1 \sim \Gamma_2 \) and \( I \) is the \( 4 \times 4 \) identity matrix. In deriving equation (6), one ignores the second and higher powers of both \( \delta E / E_M \) and \( \Gamma / E_M \) as \( E_M \gg \Delta E, \Gamma \). It is easy to check that despite the above approximation, \( S_{\delta E} \) is unitary; furthermore, equation (6) holds for any value of the ratio \( \delta E / \Gamma \) as long as \( E_M \gg \Delta E, \Gamma \). Using \( S_{\delta E} \), one readily obtains that the probability of observing an electron–electron output state becomes finite and equal to \( \frac{(\delta E)^2}{(\delta E)^2 + \Gamma^2} \), which, of course, vanishes when \( E \sim E_M \).

Before ending this section, as a brief aside, we point out that when one electron and one hole scatter at sites 1 and 2, respectively, for \( \delta E \ll \Gamma \), the incoming state \( c_1^\dagger d_2^\dagger |0\rangle \) evolves to

\[
\frac{1}{2}(-c_1^\dagger c_2^\dagger - ic_1^\dagger d_1^\dagger + ic_2^\dagger d_2^\dagger - d_1^\dagger d_2^\dagger)|0\rangle,
\]

implying the interferometric vanishing of the probability of one outgoing electron and one outgoing hole in separate leads. We will not delve into this
case further, but now proceed to discuss the signatures of the Hanbury–Brown interferometry in the case of two incident electrons.

4. Spectral density of current fluctuations: a signature of fermion antibunching in pseudospin space

In the previous section, we described the scattering as a process where one sends particles one by one through the leads at specific times. However, in practice, rather than controlling times, one can control the energies \( \epsilon_1 \) and \( \epsilon_2 \) of the particles in their respective leads so as to make them behave indistinguishably when \( \epsilon_1 \sim \epsilon_2 \). Then, the standard way to observe the predicted fermion anti-bunching is through a measurement of the correlations between the currents in leads 1 and 2. The current in lead \( j \) may be written as

\[
I_j(t) = \frac{e}{h} \sum_{\epsilon, \epsilon'} e^{i(\epsilon - \epsilon')t} \{ a_j^\dagger(\epsilon) a_j(\epsilon') - b_j^\dagger(\epsilon) b_j(\epsilon') \},
\]

(7)

where \( a_j \) and \( b_j \) denote the incoming and outgoing particles, \( \epsilon \) and \( \epsilon' \) are the energies of the particles and \( \nu \) is the density of states of the incoming electrons. The spectral density of the current fluctuations \( \delta I_j = I_j - \langle I_j \rangle \) between the leads at zero frequency is

\[
P_{ij} = \lim_{T \to \infty} \frac{h \nu}{T} \int_0^T \text{d}t \text{Re} \langle \delta I_1(t) \delta I_2(0) \rangle.
\]

(8)

Using \( S_{\delta E} \) of equation (6) and considering an incoming two-electron state \( c_1^\dagger(\epsilon_1) c_2^\dagger(\epsilon_2) |0\rangle \), where \( c_j^\dagger(\epsilon_j) \) denotes an electron of energy \( \epsilon_j \) in lead \( j \), one finds that

\[
P_{ij} = \frac{e^2}{h \nu} \frac{\Gamma^2}{\{(\delta E)^2 + \Gamma^2\}^2} (2 \delta E)^2 (\delta E^2 - \Gamma^2) \delta_{\epsilon_1, \epsilon_2},
\]

(9)

where \( \delta_{\epsilon_1, \epsilon_2} \) is the Kronecker delta function. Note that if the incident electrons are distinguishable, i.e. \( \epsilon_1 \neq \epsilon_2 \), then, as expected, \( P_{ij} = 0 \) since for an electron exiting one lead there could equally well be an electron or a hole exiting the other lead. If, instead, \( \epsilon_1 = \epsilon_2 \) (i.e. the particles are indistinguishable), then for \( |\delta E| < |\Gamma| \), the domination of the electron–hole final state (as in equation (5)) makes \( P_{ij} < 0 \), which allows us to detect the predicted ‘anti-bunching’ in pseudospin space. For \( |\delta E| > |\Gamma| \), a process of amplitude \( \Gamma^2 \delta E \) in which only one of the electrons scatters, while the other remains in its lead, dominates; Fermi statistics now makes the electrons anti-bunch spatially (the more conventional antibunching [2, 21]), thus contributing to a positive \( P_{ij} \). As in [8], our results are not inconsistent with those of Bolech and Demler [22], since their results apply when the energy of the incoming electrons is much higher than \( E_M \).

5. Scattering matrix in the spinfull case

Our analysis has so far been confined to the spinless model investigated in [7, 8], whereas for the promising implementations [12–15], the Majorana modes should involve superpositions of operators of different spins. For example, for a realization in a ferromagnet–s-wave
superconductor–ferromagnet structure on a quantum spin Hall edge [13], one has

\[
\gamma_1 = \frac{1}{\sqrt{2}} (c_{1,\uparrow} - ic_{1,\downarrow} + ic_{1,\downarrow}^\dagger + c_{1,\uparrow}^\dagger),
\]

\[
\gamma_2 = \frac{1}{\sqrt{2}} (c_{2,\uparrow} + ic_{2,\downarrow} - ic_{2,\downarrow}^\dagger + c_{2,\uparrow}^\dagger),
\]

where \(c_{i,\sigma}\) creates an electron with spin \(\sigma\) in lead \(i\). Defining the spin states \(|\pm y\rangle = \frac{1}{\sqrt{2}}(|\uparrow\rangle \pm i|\downarrow\rangle\) and using the basis \(|e_{1,\pm y}\rangle, |e_{2,\pm y}\rangle, |h_{1,\pm y}\rangle, |h_{2,\pm y}\rangle\), one obtains the scattering matrix found to be

\[
S_{\text{spinfull}} = \begin{pmatrix} I & 0 \\ 0 & S \end{pmatrix},
\]

where in (11), \(I\) and \(0\) are the \(4 \times 4\) identity and null matrices, while \(S\) is the scattering matrix given by equation (3).

When one uses \(S_{\text{spinfull}}\) to study the scattering of the incident state \(c_{1,-y}c_{2,\pm y}|0\rangle\), one only needs the lower-right \(4 \times 4\) block of \(S_{\text{spinfull}}\). Thus, precisely the same electron–hole output state as that in equation (5) is obtained, apart from the fact that now the spin indices \(-y\) and \(+y\) are pinned to sites 1 and 2, respectively. Thus, by choosing the spin polarizations of the incoming electrons pertinently, one can observe all the effects described until now. This should be possible in a variety of systems as Majorana modes of the form given by equation (10) are quite generic, e.g. also realizable in semiconductor–superconductor–magnet structures [15].

6. Entanglement of distant electron spins

We now propose a protocol for generating of entanglement between spins of well-separated particles incoming at site 1 and at site 2. For this purpose, we choose the realization of Majorana fermions given by equation (10) and make two electrons with parallel spins in the \(\uparrow\) direction come in simultaneously, i.e. choose the initial state \(c_{1,\uparrow}^\dagger c_{2,\uparrow}^\dagger|0\rangle\). Then, using \(S_{\text{spinfull}}\), one obtains

\[
c_{1,\uparrow}^\dagger c_{2,\uparrow}^\dagger|0\rangle \xrightarrow{\text{MBS}} \frac{1}{4}(c_{1,\uparrow}^\dagger c_{2,\uparrow}^\dagger - c_{1,\downarrow}^\dagger c_{2,\downarrow}^\dagger + 2c_{1,\uparrow}^\dagger c_{2,\downarrow}^\dagger + \cdots)|0\rangle,
\]

where \(\cdots\) denotes terms such as \(c_{1,\sigma}^\dagger c_{1,\sigma}^\dagger\), \(c_{2,\sigma}^\dagger c_{2,\sigma}^\dagger\), \(c_{1,\uparrow}^\dagger d_{k,\alpha}^\dagger\), \(c_{1,\downarrow}^\dagger d_{k,\alpha}^\dagger\), \(c_{2,\uparrow}^\dagger d_{k,\alpha}^\dagger\), \(c_{2,\downarrow}^\dagger d_{k,\alpha}^\dagger\), and \(d_{k,\alpha}^\dagger d_{k,\alpha}^\dagger\), which are not relevant to our discussion. Equation (12) implies that when two outgoing electrons are obtained in leads 1 and 2 separately, their state is \(|\xi\rangle_{12} = \frac{1}{\sqrt{2}}(|\uparrow\rangle_1 |\uparrow\rangle_2 - |\downarrow\rangle_1 |\downarrow\rangle_2 + 2|\uparrow\rangle_1 |\downarrow\rangle_2\), where, as is usually done [5, 21], one uses the lead label to label the electron. \(|\xi\rangle_{12}\) is an entangled state of the spins of electrons 1 and 2, with the amount of entanglement (as quantified by the von Neumann entropy of one of the particles [24]) being 0.19 ebit. Although the entanglement is not very high, \(|\xi\rangle_{12}\) is a pure state and hence of value in quantum information, as its entanglement can be concentrated without loss by local means [24]. Moreover, the probability of obtaining two outgoing electrons in separate leads (i.e. \(|\xi\rangle_{12}\) is rather high, namely 3/8. At the expense of decreasing this probability, one may improve the degree of entanglement of the generated state by tuning the polarizations of the incoming electrons. For instance, if the incoming state is \((\frac{1}{\sqrt{3}} c_{1,\uparrow}^\dagger + \frac{1}{\sqrt{3}} c_{1,\downarrow}^\dagger + \frac{1}{\sqrt{3}} c_{1,\downarrow}^\dagger)(\frac{1}{\sqrt{3}} c_{2,\uparrow}^\dagger + \frac{1}{\sqrt{3}} c_{2,\downarrow}^\dagger)|0\rangle\), one obtains an output state of entanglement 0.75 ebits while the probability of generating this state becomes 0.055. The spin entanglement of the outgoing electrons could be measured by passing them through separate spin filters as in [23].
Unlike the entanglement generation scheme of [5], here particles polarized parallel to each other suffice to generate entanglement. Importantly, in our protocol, particles at a distance from each other can be made entangled; this may avoid the decoherence arising necessarily from the transport needed to separate the particles after a local entangling mechanism. In addition, the distance between the entangled particles can be enhanced by putting $n$ copies of our setup in series with leads connecting the end of one copy to the beginning of another. The probability of obtaining the state $|\xi\rangle_{12}$ in the leftmost and rightmost leads will then be $(3/8)^n$.

One might think that in analogy with [21], perhaps it is possible also to detect entanglement between distant electronic spins by injecting to the opposite ends of our one-dimensional device. As a maximally entangled state of two spins can be written in any basis, let us consider the $|\pm y\rangle$ basis for spins. In this basis, two of the maximally entangled states can be written as $|\psi^\pm\rangle_{12} = (c_{1,+,y}^\dagger c_{2,-y}^\dagger \pm c_{1,-y}^\dagger c_{2,+,y}^\dagger) |0\rangle$. The detection of entanglement at a normal 50:50 beam splitter relies crucially on both the incoming states $c_{1,+,y}^\dagger c_{2,-y}^\dagger |0\rangle$ and $c_{1,-y}^\dagger c_{2,+,y}^\dagger |0\rangle$ evolving at the beam splitter and interference (i.e. cancellation/addition) between the terms resulting from the evolution of each of the above two states. Only as a result of these cancellations/additions do the bunching/anti-bunching effects evidencing entanglement arise. However, here the term $c_{1,+,y}^\dagger c_{2,-y}^\dagger |0\rangle$ does not even evolve under the action of $S_{\text{spinfull}}$ so that interferences are impossible. Thus although our device can generate spin entanglement it cannot detect spin entanglement.

7. A few condensed matter settings

One simple setting where the non-local two-particle interferometry and the entanglement generation between distant electrons from MBS may be observed can be engineered with magnet–superconductor–magnet junctions deposited on the edge of a 2D quantum spin Hall insulator [8, 13]. Just as in [8], one can observe these effects when the Majorana modes are separated by a distance $d$ of several micrometers at temperatures of the order of 10 mK. For this setting, the explicit form of the Majorana operators is exactly the same as that in equation (10) [13]. Interestingly, strong spin–orbit coupled quantum wires in proximity to ferromagnets and superconductors also support the realization of MBS [14, 15] given in equation (10) [15]. As in previous proposals [13, 14, 22], also in these settings, two STM tips could act as leads 1 and 2 to observe the non-local two-particle Hanbury–Brown–Twiss interferometry. For the entanglement generation, instead, it will be more useful to have synchronized electron pumps [25] feeding in the incoming electrons. In addition, the filtering of the desired state $|\xi\rangle_{12}$ can be achieved by pumps capturing exactly one outgoing electron from each Majorana bound state.

8. Conclusions

In this paper, we have shown that a 1D device with two MBS at its ends yields a Hanbury–Brown–Twiss effect in the internal electron–hole pseudospin space, which may be detected in realistic condensed matter settings through current–current correlations. This is a departure from all the known multi-particle interference effects that have manifested themselves in spatial bunching and antibunching or spin–spin correlations. Fundamentally, it can be regarded as a manifestation of the quantum indistinguishability between electronic annihilation and hole creation evidencing the presence of MBS. The same settings may also be used to...
engineer a non-local entangler of distant electronic spins, which may enable one to circumvent the decoherence arising from the transport needed to separate entangled particles.

Acknowledgments

We thank C Chamon, R Egger, R Jackiw, M Pepper, S Y Pi, G W Semenoff and S Tewari for fruitful discussions. SB (PS) thanks the University of Perugia (University College London) for hospitality and partial support. SB thanks the UK EPSRC, the Royal Society and the Wolfson Foundation.

References

[1] Hong C K, Ou Z Y and Mandel L 1987 Phys. Rev. Lett. 59 2044
[2] Liu R C et al 1998 Nature 391 263–5
[3] Jeltes T et al 2007 Nature 445 402
[4] Shih Y H and Alley C O 1988 Phys. Rev. Lett. 61 2921–4
[5] Bose S and Home D 2002 Phys. Rev. Lett. 88 050401
[6] Fu L and Kane C L 2009 Phys. Rev. Lett. 102 216403
    Akhmerov A R, Nilsson J and Beenakker C W J 2009 Phys. Rev. Lett. 102 216404
[7] Semenoff G W and Sodano P 2007 J. Phys. B: At. Mol. Opt. Phys. 40 1479
[8] Nilsson J, Akhmerov A R and Beenakker C W J 2008 Phys. Rev. Lett. 101 120403
[9] Fu L 2010 Phys. Rev. Lett. 104 056402
[10] Kitaev A Yu 2001 Phys.—Usp. 44 131
[11] Tewari S et al 2007 Phys. Rev. Lett. 98 010506
[12] Fu L and Kane C L 2008 Phys. Rev. Lett. 100 096407
[13] Shivamoggi V, Refael G and Moore J E 2010 Phys. Rev. B 82 041405
[14] Sau J D et al 2010 Phys. Rev. B 82 214509
    Sau J D et al 2010 Phys. Rev. Lett. 104 040502
[15] Oreg Y, Refael G and von Oppen F 2010 Phys. Rev. Lett. 105 177002
[16] Chamon C et al 2010 Phys. Rev. B 81 224515
[17] Kitaev A Yu 2003 Ann. Phys., NY 303 1
    Nayak C et al 2008 Rev. Mod. Phys. 80 1083
[18] Datta S 1995 Electronic Transport in Mesoscopic Systems (Cambridge: Cambridge University Press) pp 145–57
    Fisher D S and Lee P A 1981 Phys. Rev. B 23 6851
[19] Ghur T, Møller-Groeling A and Weidenmuller H 1998 Phys. Rep. 299 189–425
[20] Büttiker M 1990 Phys. Rev. Lett. 65 2901
[21] Burkard G, Loss D and Sukhorukov E V 2000 Phys. Rev. B 61 R16303
[22] Bolech C J and Demler E 2007 Phys. Rev. Lett. 98 237002
[23] Sauret O, Martin T and Feinberg D 2005 Phys. Rev. B 72 024544
[24] Bennett C H et al 1996 Phys. Rev. A 53 2046
[25] Blumenthal M D et al 2007 Nature Phys. 3 343
    Wright S J et al 2009 Phys. Rev. B 80 113303

New Journal of Physics 13 (2011) 085002 (http://www.njp.org/)