

Super-Poissonian noise from quasiparticle poisoning in electron transport through a pair of Majorana bound states

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Topological qubits based on non-Abelian Majorana bound states (MBSs) are protected by parity which is challenged by quasiparticle poisoning (QPP). In this work, we show how QPP affects transport through a pair of coupled MBSs weakly connected to two current leads, using an open system approach and full counting statistics. We find that the correct low-energy physics requires one to include next to leading order tunneling events in the coupling to the leads. In particular, our results show that QPP causes super-Poissonian local shot noise with a Fano factor that diverges with decreasing bias voltage, while the current and the nonlocal (cross) noise are only little affected. We explain that these features are a direct consequence of the nature of MBSs being their own antiparticle making noise measurements a viable tool to search for MBSs in the presence of QPP.

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Majorana bound states (MBSs) are predicted to be zero-energy excitations at the boundary of topological superconductors [1–9] with the defining property that its creation operator $\gamma^\dagger$ is identical to its annihilation operator γ , but their existence has not been conclusively confirmed by experiments yet, even though they should be topologically protected and hence robust to disorder.

These non-Abelian quasiparticles have been advertised as building blocks for topological quantum computing [3,10–15]. The superior power of such qubits lies in the nonlocal encoding of quantum information. These qubits, however, are prone to parity changes, i.e., number of particles in the superconductor being odd or even. Parity switching events are therefore detrimental for topological quantum computing. An exception is their controlled use in braiding processes of MBSs [16]. However, uncontrolled (i.e., random) parity changes appear naturally in open quantum systems and are also referred to as quasiparticle poisoning (QPP) [17–22].

The QPP mechanism is relevant not only for systems hosting MBSs, but also for superconducting (nontopological) qubits [23–26] and for quantum dot based topological superconductor devices [27].

In this Letter, we consider the transport properties of a pair of MBSs comprising a nonlocal fermion state as implemented, e.g., in a Majorana wire and attached to two leads; see Fig. 1(a). We apply an open systems approach for the

calculation of the cumulant generating function (CGF) leading to average current and noise properties (and in principle to higher-order cumulants) in the presence of dissipation. In particular, we study the effects of QPP.

Transport through such and related setups has been extensively studied in the absence of QPP [28–47] theoretically and experimentally [48–52]. In the low voltage bias regime, below the hybridization energy of the two MBSs, and at weak coupling, the cross noise properties have been shown to be perfectly correlated, hinting at the splitting of a Cooper pair via two MBSs as the dominant transport process [29,35]. The local Fano factor is therefore unity. This physics has been attributed to the coupling of the tunneling electrons to single MBSs. In this work, we show that, in the same low bias regime, the inclusion of QPP can lead to a dramatic change of these results, in particular, to a diverging super-Poissonian local Fano factor for the shot noise. We attribute this physics to the possible nonequilibrium population of the higher nonlocal fermion state of the MBS pair driven by the QPP rate and of the fact that a MBS couples to electrons and holes with the same strength by virtue of $\gamma = \gamma^\dagger$.

All calculations are performed analytically, yielding expressions for the transport quantities that are discussed in the results sections. We also find that the low bias regime cannot be described by lowest order tunneling (sequential tunneling, ST) events alone. Second-order tunneling (SOT) events (Andreev tunneling and electron cotunneling) we treat by deriving an effective Hamiltonian.

Model Hamiltonian. A pair of coupled MBSs, each residing at an end of a proximitized quantum wire put to ground is tunnel coupled to two normal metallic leads; see Fig. 1(a).

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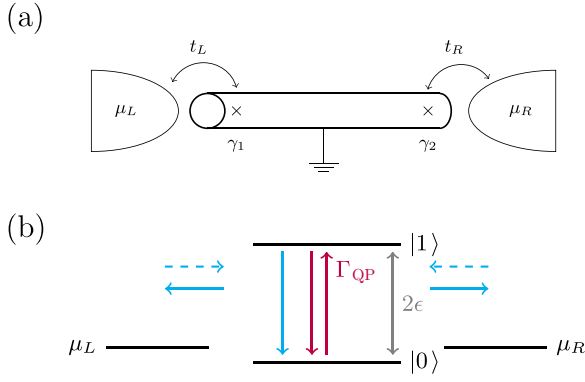


FIG. 1. (a) Sketch of the transport setup. (b) Energy diagram with possible ST at low chemical potentials μ_L and μ_R . Deexciting the MBSs system (vertical blue arrow) can be associated with both tunneling out from the wire, via normal tunneling (blue horizontal arrows), or into the wire, via anomalous tunneling (blue horizontal vertical arrows). QPP (dark red arrows) switches the parity state with associated energy cost 2ϵ without any tunneling processes. SOT processes are not shown here, as they do not change the parity. However, they can contribute to current and noise.

The Hamiltonian for the system is given by (we set $\hbar = 1$) $H = H_\gamma + H_B + H_T$, where

$$H_\gamma = i\epsilon\gamma_1\gamma_2 = -\epsilon(1 - 2c^\dagger c) \quad (1)$$

describes the two coupled MBSs with self-adjoint operators γ_1, γ_2 and hybridization energy ϵ and where we have used fermion creation (annihilation) operators $c^\dagger$ (c) to rewrite the Majorana operators $\gamma_1 = c^\dagger + c$ and $\gamma_2 = i(c^\dagger - c)$. The leads are modeled by the bath Hamiltonian

$$H_B = iv_F \sum_{i=L,R} \int_{-\infty}^{\infty} dx_i \psi_i^\dagger(x_i) \partial_{x_i} \psi_i(x_i), \quad (2)$$

where $\psi_i(x)$ [$\psi_i^\dagger(x)$] annihilate (create) chiral electrons at position x in lead $i = L, R$ moving with the Fermi velocity v_F . The MBSs are tunnel coupled to leads by the tunneling Hamiltonian

$$H_T = t_L \psi_L \gamma_1 + t_R \psi_R \gamma_2 + \text{H.c.} \\ = t_L \psi_L c^\dagger + it_R \psi_R c^\dagger + t_L \psi_L c - it_R \psi_R c + \text{H.c.}, \quad (3)$$

where $t_{R(L)}$ are amplitudes for tunneling of an electron at $x = 0$ from lead $R(L)$ to the corresponding MBS; see Fig. 1(a). The first two terms of the second row correspond to normal tunneling and the second two terms of the second row correspond to anomalous tunneling. The normal tunneling conserves charge in the system, while the anomalous tunneling only conserves parity. In the Supplemental Material [53], we consider a more general tunneling Hamiltonian, allowing for nonlocal tunneling amplitudes.

Master equation and full counting statistics. To calculate transport properties and to include dissipative processes like tunneling to the leads and QPP events we employ a master equation approach and full counting statistics. We are mainly interested in the low-bias regime $eV \ll \epsilon$ where the coupling of electrons and holes to MBSs leads to peculiar transport properties like enhanced crossed Andreev reflection in the

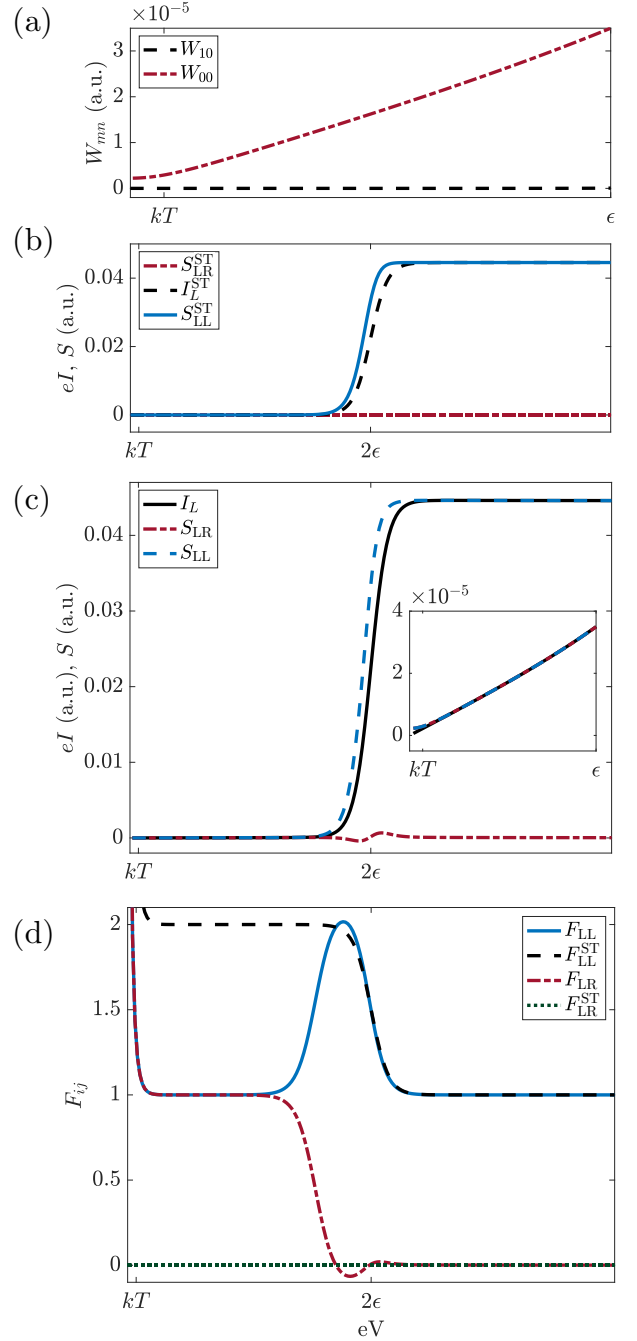


FIG. 2. (a) ST and SOT rates for a symmetrically applied bias voltage, $\mu_L = \mu_R = eV$, shown for small values of $eV < \epsilon$. Current and noise for the system without QPP (b) for ST only and (c) for ST and SOT. (d) Corresponding Fano factors. Parameter values used are $k_B T = 0.07\epsilon$ and $\Gamma_i = 7 \times 10^{-3}\epsilon$ ($i \in \{L, R\}$).

absence of QPP [29,35]. ST events from the ground state $|0\rangle$ of H_γ with $c|0\rangle = 0$ to the occupied (excited) state $|1\rangle$ with $c^\dagger c|1\rangle = |1\rangle$ and energy ϵ are strongly suppressed (for $eV, k_B T \ll \epsilon$). SOT events, on the other hand, preserve the parity $\langle c^\dagger c \rangle$ and are dominant at small energies; see Fig. 2(a). To set up the full counting statistics, we introduce counting fields $\chi = \{\chi_i\}$ counting particles in lead $i = L, R$. To include SOT processes in the master equation, we derive an

effective Hamiltonian [53,54] to second order in H_T , leading to $H^{\text{eff}} = H + H^{\text{SOT}}$, with

$$H^{\text{SOT}} = |0\rangle\langle 0| \sum_{a,b=1}^4 (-1)^{a+1} \alpha_a \alpha_b F_a \frac{1}{H_B - E_0 + 2\epsilon} F_b + |1\rangle\langle 1| \sum_{a,b=1}^4 (-1)^{b+1} \alpha_a \alpha_b F_a \frac{1}{H_B - E_0 - 2\epsilon} F_b, \quad (4)$$

where a and b run over the four terms in the original tunneling Hamiltonian H_T , $\alpha_{a,b} \in \{t_L, t_R, t_L^*, t_R^*\}$ and $F_{a,b} \in \{\psi_L, \psi_L^\dagger, \psi_R, \psi_R^\dagger\}$ are the corresponding bath field operators, and E_0 is the energy of the state of the bath.

The counting field dependent density matrix of the total system $\rho_{\text{tot}}(\chi, t)$ satisfies a generalized von Neumann equation $\dot{\rho}_{\text{tot}}(\chi, t) = -i\mathcal{L}_{H^{\text{eff}}} \rho_{\text{tot}}(\chi, t)$ with the Liouvillian superoperator $\mathcal{L}_{H^{\text{eff}}} X = [H^{\text{eff}}(\chi/2)X - XH^{\text{eff}}(-\chi/2)]$. Here, we have introduced $H^{\text{eff}}(\chi) = \exp(-i \sum_i \chi_i N_i) H^{\text{eff}} \exp(i \sum_i \chi_i N_i) = H_B + H_\gamma + H_T(\chi) + H^{\text{SOT}}(\chi)$ with N_i the number operator in lead i . Following standard steps [55,56] and assuming $\max\{eV, k_B T\}, 2\epsilon \gg \Gamma_i$, we derive a master equation for the reduced density matrix $\rho(\chi, t) = \text{Tr}_B \rho_{\text{tot}}(\chi, t)$ to lowest order in $H_T(\chi) + H^{\text{SOT}}(\chi)$

$$\dot{\rho}(\chi, t) = -\mathcal{L}(\chi) \rho(\chi, t), \quad (5)$$

where Tr_B means taking the trace over the bath degrees of freedom. We observe that the nondiagonal elements of ρ decouple, so that we can vectorize its diagonal entries $\langle n | \rho(\chi, t) | n \rangle \equiv \rho_{nn}(\chi, t)$ and the corresponding Liouvillian matrix $\mathcal{L}$ can be decomposed as $\mathcal{L} = \mathcal{L}_{\text{ST}} + \mathcal{L}_{\text{SOT}}$, with (in the basis $[\rho_{00}(\chi, t), \rho_{11}(\chi, t)]^T$)

$$\mathcal{L}_{\text{ST}} = \begin{pmatrix} W_{10}(0) & -W_{01}(\chi) \\ -W_{10}(\chi) & W_{01}(0) \end{pmatrix} \quad (6)$$

accounting for the ST processes and

$$\mathcal{L}_{\text{SOT}} = \begin{pmatrix} -W_{00}(\chi) & 0 \\ 0 & -W_{11}(\chi) \end{pmatrix} \quad (7)$$

for the SOT processes.

The ST rates from a parity state $|n\rangle$ to a parity state $|m\rangle$ ($m \neq n$) are found to be of the form $W_{mn} = W_{mn}^L + W_{mn}^R$, with

$$W_{mn}^R(\chi) = \Gamma_R [(1 - f_R^{n,m}) e^{-i\chi_R} + f_R^{m,n} e^{i\chi_R}], \quad (8)$$

where $\Gamma_R = 2\pi \rho_R |t_R|^2$ with $\rho_R = (2\pi v_F)^{-1}$ the constant density of states in the right lead and $f_R^{n,m} = f_R(E_n - E_m)$ the Fermi function with chemical potential μ_R and temperature T and $W_{mn}^L = W_{mn}^R (R \leftrightarrow L)$. Analytical expressions for the total SOT rates $W_{00}(\chi)$ and $W_{11}(\chi)$ can be found in the Supplemental Material [53]. The rates $W_{mn}(\chi = 0)$ and $W_{nm}(\chi = 0)$, $n \neq m$, are depicted in Fig. 2(a).

The CGF is defined by $S(\chi, t) = -\ln[\text{Tr}(\rho(\chi, t))] \approx \Lambda_0(\chi)t$ [57,58], where the trace is to be taken over the small system only. The last (approximate) equality is valid for long times $t = \tau$, where the eigenvalue with the smallest real part, Λ_0 , is adiabatically connected to zero for $\chi \rightarrow 0$.

The mean current and the current noise can be found from the χ derivatives of the CGF [58], taken at $\chi = 0$

$$I_i = i \frac{e}{\tau} \frac{\partial}{\partial \chi_i} S(\chi, \tau)|_{\chi=0}, \quad (9)$$

$$S_{ij} = \frac{e^2}{\tau} \frac{\partial^2}{\partial \chi_i \partial \chi_j} S(\chi, \tau)|_{\chi=0}. \quad (10)$$

From these quantities, we calculate the Fano factors F_{ij} as [59]

$$F_{ij} = \frac{S_{ij}}{e|I_i|}. \quad (11)$$

Without QPP terms, the parity state $|1\rangle$ is never occupied at low bias voltages, so in this case there is approximately no current from ST processes. When a QPP rate is introduced, the parity is allowed to flip without any tunneling process to/from the leads taking place.

To account for such an effective QPP process, we add a phenomenological term to the Liouvillian $\mathcal{L} \rightarrow \mathcal{L}_{\text{ST}} + \mathcal{L}_{\text{SOT}} + \mathcal{L}_{\text{QP}}$. The Liouvillian $\mathcal{L}_{\text{QP}}$ changes the parity in the system and is given by (cf. Fig. 1) [60–64]

$$\mathcal{L}_{\text{QP}}[\rho] = -\Gamma_{\text{QP}}(c\mathcal{P}\rho\mathcal{P}^\dagger - \frac{1}{2}\{\rho, \mathcal{P}c^\dagger c\mathcal{P}\}) + \text{H.c.}, \quad (12)$$

with $\mathcal{P} = (-1)^{c^\dagger c}$. For our system we can write this in matrix form as

$$\mathcal{L}_{\text{QP}} = \Gamma_{\text{QP}} \begin{pmatrix} 1 & -1 \\ -1 & 1 \end{pmatrix}. \quad (13)$$

Current and noise without QPP. For concreteness, we consider a symmetric bias case $\mu_L = \mu_R = eV > 0$ where current flows from the leads to the superconductor [29,35]. The transport properties as a function of V are fundamentally different at low V ($eV \ll 2\epsilon$) and large V ($eV \gg 2\epsilon$). At small V , the probability to be in $|1\rangle$ is exponentially small at low temperatures and transport is dominated by H^{SOT} only virtually occupying the excited state [the rate is only suppressed by $(1/\epsilon)^2$]. At large V , ST processes are dominant and SOT processes are negligible. It turns out that, as an exact result, ST events alone cannot lead to cross correlations, i.e., $S_{\text{LR}}^{\text{ST}} = 0$ [cf. Fig. 2(b)]. However, SOT processes are important at small V [cf. Fig. 2(c)]. As a consequence, the ST-only result contributes wrong Fano factors at small V , whereas the result including SOT contributions agrees quantitatively with Refs. [29,35] at small and large V ; see Fig. 2(d) [65]. We note that, when $eV \simeq 2\epsilon$, our calculation is not valid since then the SOT contributions become ill defined. The uprise of F_{ij} for $eV \lesssim k_B T$ is a consequence of dominant thermal noise contributions.

Current and noise with QPP. In the presence of QPP, the two parity states $|0\rangle$ and $|1\rangle$ become statistically occupied with a switching rate Γ_{QP} ; see Eq. (13). We observe that the average current I_L is only weakly affected by QPP and, while the cross noise S_{LR} is largely unaffected, the local noise S_{LL} is significantly offset to larger values in the presence of QPP [cf. Fig. 3(a)]. As a consequence, the local Fano factor F_{LL} becomes much larger than unity and shows a diverging tendency towards small voltages [cf. Fig. 3(b)]. Note that S_{LL} is significantly enhanced compared to the equilibrium noise when $eV \rightarrow 0$ (compare with S_{LR} and F_{LR} which are largely unaffected by QPP) underlining that QPP is a nonequilibrium

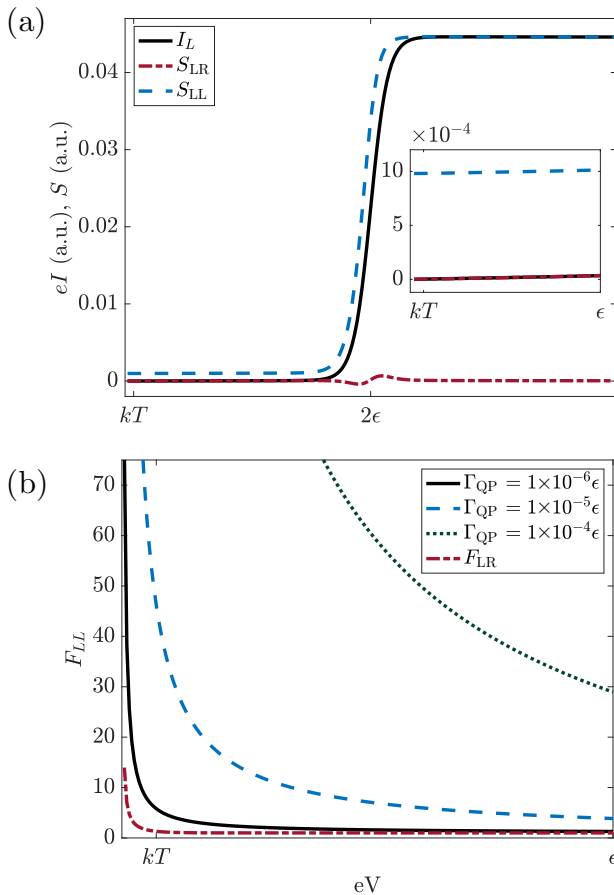


FIG. 3. (a) Current and noise as a function of bias voltage for the model including QPP with $\Gamma_{QP} = 1 \times 10^{-4}\epsilon$ (including both ST and SOT processes). (b) Local noise Fano factor for different values of Γ_{QP} . The cross noise Fano factor F_{LR} is unaffected by the finite QPP. Parameters used are $k_B T = 0.07\epsilon$ and $\Gamma_i = 7 \times 10^{-3}\epsilon$ ($i \in \{L, R\}$).

effect. These features can be explained considering that QPP induces a significant weight for the system to be in the excited state $|1\rangle$. This opens up the possibility for strong ST events with rate W_{01} even at small V , bringing the system back to the ground state $|0\rangle$. As we can deduce directly from Eq. (8) for $\chi = 0$, the ST rate has a normal part $\propto [1 - f_L(2\epsilon)]$ tunneling an electron to the lead and an anomalous part $\propto f_L(-2\epsilon)$

creating a hole in the lead (thereby forming a Cooper pair in the superconducting condensate). Since these processes involve the coupling to a single MBS, the two contributions have the same strength Γ_L (and similar for the right lead). This MBS feature explains the strong S_{LL} (the two parts are added) in the absence of an average current I_L (the two parts are subtracted) induced by these strong ST events. Note that, for $eV, k_B T \ll 2\epsilon$, $[1 - f_{L,R}(2\epsilon)] \approx f_{L,R}(-2\epsilon) \approx 1$ so that the average ST current is essentially zero and the main contribution to I_L originates from crossed Andreev reflection processes (SOT contributions) [29,35] with $W_{mn} \propto \Gamma_L \Gamma_R eV / \epsilon^2$ for $k_B T \ll eV$ [53]. In the same limit, we obtain for the Fano factor [53] $F_{LL} \approx 1 + \bar{\rho}_{11} 2\pi \epsilon^2 / (\Gamma_R eV)$ with $\bar{\rho}_{11}$ the probability to be in state $|1\rangle$ [satisfying $\mathcal{L}(\chi = 0)\bar{\rho} = 0$]. This explains the power-law increase of F_{LL} with decreasing V observed in Fig. 3(b) as well as the fact that QPP only very weakly affects the average current and nonlocal properties (induced by SOT processes). We note that, in the Majorana-based setups studied in [60,66,67], Γ_{QP} leads to a suppression of shot noise.

In conclusion, we showed that QPP leads to dramatic effects on current noise properties of a coupled pair of MBSs weakly tunnel coupled to current leads including ST and SOT events in an open quantum system approach. Particular emphasis was put on the regime of small bias voltage (below the hybridization energy of the two MBSs) applied between the leads and the grounded superconductor. The local Fano factor was shown to grow with decreasing bias voltage in a power-law fashion (reaching values much larger than one), whereas the average current and nonlocal current fluctuations are largely unaffected by QPP. This striking feature is based on the fact that average current and nonlocal current fluctuations are dominated by (small) SOT events, whereas the local noise is dominated by (larger) ST processes. We explained how these QPP-induced noise effects can be traced back to the coupling of each lead to a single MBS. This allows one to test characteristic MBS features via electron transport in the presence of QPP.

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