

Manipulation of Majorana Modes in a Double Quantum Dot

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(Dated: February 7, 2019)

Luis We'll call them Majorana "zero modes" instead of "Majorana fermions", which relates more to elementary particles.

Majorana zero modes (MZMs) appearing at the edges of topological superconducting wires are a promising platform for fault-tolerant quantum computation. Novel proposals use MZMs tunneling inside quantum dots (QDs) to implement quantum architectures, because today's precise experimental control over the QD parameters offers the unique possibility of manipulating the Majorana modes inside multi-dot systems. The simplest case where Majorana manipulation is possible is in a double quantum dot (DQD). So far, no complete analysis of this basic structure has been done. This project fills this gap by realizing an exact quantum transport study of the effects of coupling a Majorana mode with a DQD. By tuning the model parameters we show that it is possible to control the localization of the Majorana signature in the DQD.

I. INTRODUCTION

→ Majorana zero modes in condensed matter systems: they have been found, several papers have been written about it and there has been much progress in distinguishing them from other sources of zero-bias peaks.

In the last few decades the interest in the so called Majorana fermions has been increasing. The particle proposed by the physicist Ettore Majorana as the real field solution of the Dirac equation describes a fermion which is its own antiparticle, hence it has no charge nor mass. To the date no fundamental particle with these characteristics has been found. However, theoretical research predicts that Majorana Fermions emerge as quasi-particles at the boundary of certain topological superconductors. Kitaev¹ Recently, the new technological innovations allowed the observation of Majorana signatures in different topological materials.²⁻⁵

→ What's next? Manipulation and braiding !

Despite the positive experimental results, their is still certain skepticism about the existence of Majorana Fermions. One of the reasons is that some properties of Majorana quasiparticles like the expected non-abelian statistics have not been measure. This property is of especial interest due to its promising applications in topological quantum computing. The braiding protocol based on Majorana's non-abelian statistics is the key to fault-tolerant quantum computation.^{6,7}

→ MZMs in quantum dots can co-exist with Kondo peaks.

A promising method to detect Majorana modes consists in attaching a quantum dot (QD) to the edges of a Majorana chain in the topological phase and executing transport measurements through the QD.⁸ The Majorana mode at the end of the chain then leaks inside the QD⁹ which produces a zero-bias conductance peak of half a quanta $\frac{e^2}{2h}$ through the dot. This is a Majorana signature which produces half of the expected peak by a regular fermion. Recently, experiments including hybrid Majorana-QD systems have been performed.¹⁰ In addition,

the similarity of this phenomenon with the Kondo effect,^{11,12} where the zero-bias conductance peak takes $\frac{e^2}{h}$, motivated the study of combined Kondo-Majorana physics in this system.^{13,14} This project revealed the existence of a region of parameters where both, Kondo and Majorana physics, coexist.

This idea has turned on new lights into the design of quantum architectures,^{15,16} because today's precise experimental control over the parameters of QDs - energy levels, tunneling couplings, etc. - offers the unique possibility of manipulating the Majorana modes inside multi-dot systems. The simplest case where Majorana manipulation is possible is in a double quantum dot. So far, no complete analysis of this simple case has been done. The goal of this project is to fill this gap by realizing a full quantum transport study of the effects of coupling a Majorana mode with a double quantum dot. By tuning the QD gate voltages and the Majorana couplings we will be able to probe the mobility of the Majorana modes through the dots.

Luis Ok. Ref. 16 is PRB of scalable designs: PRB **95** 235305 (2017).

→ Here's what we did: quantum tunneling of a MZM into a double dot shows several possibilities for manipulation of MZM

Luis In this paper, we ... We considered both interacting and non-interacting cases. For interacting systems we used a obtained the exact transport description . On non-interacting models we used a NRG approach. We found that in symmetric couplings In the non-interacting case, we confirmed that shifting the QD's gate voltage induces the Majorana to tunnel only to the other dot. In addition, an indirect coupling of the second dot could cause destructive interference with the Majorana signature. In the interacting case, the NRG simulations confirmed these results and showed that other interacting effects - Kondo effect and RKKY interactions¹⁷⁻¹⁹ - could coexist with the Majorana signatures. On the other hand, when only one QD is coupled to the leads and the other Dot is attached to the QD, the Kondo effect is an-



FIG. 1: DQD-Majorana set-up. Solid lines: standard coupling. Dashed lines: Majorana spin- \downarrow effective couplings (6). The atomic energy levels appear inside each QD. Red dashed horizontal lines represent the Fermi level.

nihilated due to the destructive interference generated by extra dot.²⁰ Our study includes how the Majorana mode interacts with these two effects.

II. MODEL AND METHODS

We consider the setup shown in Figure 1 in which a Majorana mode at the edge of Topological Superconductor(TS) is coupled to a double quantum dot (DQD), which is attached to a single metallic lead. The Hamiltonian of this system can be partitioned in four terms: the DQD Hamiltonian H_{DQD} , the Lead Hamiltonian H_{Lead} , the DQD-lead interaction $H_{DQD-Lead}$ and the coupling between the DQD and the Majorana mode H_{M-DQDs} and

$$H = H_{DQD} + H_{Lead} + H_{DQD-Lead} + H_{M-DQD} \quad (1)$$

The interacting Anderson Model describes the DQD-lead system

$$H_{DQD} = \sum_{i \in \{1,2\}} \sum_{\sigma \in \{\downarrow, \uparrow\}} \left(\epsilon_{di} + \frac{U_i}{2} \right) \hat{n}_{i\sigma} + \frac{U_i}{2} (\sum_{\sigma} \hat{n}_{i\sigma} - 1)^2 + \sum_{\sigma \in \{\downarrow, \uparrow\}} t_{dots} (d_{1\sigma}^\dagger d_{2\sigma} + d_{2\sigma}^\dagger d_{1\sigma}), \quad (2)$$

and

$$H_{Lead} = \sum_{\mathbf{k}\sigma} \epsilon_{\mathbf{k}} c_{\mathbf{k}\sigma}^\dagger c_{\mathbf{k}\sigma} \quad (3)$$

$$H_{DQD-Lead} = \sum_{\mathbf{k}\sigma} \sum_{i \in \{1,2\}} V_{i\mathbf{k}} c_{\mathbf{k}\sigma}^\dagger d_{i\sigma} + V_{i\mathbf{k}}^* d_{i\sigma}^\dagger c_{\mathbf{k}\sigma}, \quad (4)$$

where ϵ_{di} is the energy level of dot i , U_i is the Coulomb repulsion and t_{dots} is the coupling parameter between

both QDs. The operator $d_{i\sigma}^\dagger$ creates a particle in dot i with spin σ and $\hat{n}_{i\sigma} := d_{i\sigma}^\dagger d_{i\sigma}$ is the particle number operator of state i . $c_{\mathbf{k}\sigma}^\dagger$ is the creation operator a particle with momentum \mathbf{k} and spin σ in the lead. $\epsilon_{\mathbf{k}l}$ is the corresponding energy and $V_i(\mathbf{k})$ describes the tunneling coupling between the lead and dot i .

The Majorana modes are modeled as a superposition of the creation and annihilation operators of a spin \downarrow particle f_\downarrow

$$\gamma_1 := \frac{1}{\sqrt{2}} (f_\downarrow^\dagger + f_\downarrow), \gamma_2 := \frac{i}{\sqrt{2}} (f_\downarrow^\dagger - f_\downarrow). \quad (5)$$

This makes possible to define an effective coupling between the Majorana Mode and the DQD by attaching γ_1 with the spin- \downarrow channel in the QDs

$$H_{M-DQD} = \sum_{i=1}^2 t_i (d_{i\downarrow}^\dagger \gamma_1 + \gamma_1 d_{i\downarrow}) + \epsilon_M \gamma_1 \gamma_2. \quad (6)$$

where t_i is the coupling parameter between the Majorana mode and QD i . ϵ_m is the coupling energy between both Majorana modes.

A. Methods

B. Non-interacting system

Using Zubarev's ballistic transport approach²¹, we derived an exact expression for the Green functions associated to both quantum dot operators ($G_{d_1 d_1^\dagger}(\omega), G_{d_2 d_2^\dagger}(\omega)$). The detailed procedure is included in Appendix A. The transport equations define a 9×9 linear system where the Hamiltonian parameters ($t_1, t_2, \epsilon_1 \dots$) and the energy ω are taken as algebraic variables. The solution of this type of equations is a polynomial fraction of the same degree which makes difficult to provide an exact solution using analytical or numerical methods. To bypass this problem, we associated this transport system to a flow graph and executed a Graph-Gauss-Jordan elimination process²². This method proved to be efficient to solve complex transport systems since the graph structure allows us to identify minimum cutting points and create an algorithmic representation of the Green function.

At the end, we obtained the following analytical expression

$$G_{d_{1\downarrow}, d_{1\downarrow}^\dagger}(\omega) = \frac{1}{\omega - \epsilon_{DQD}^+ - \frac{\|T_+\|^2}{\omega - \epsilon_{M2} - \frac{\|T_-\|^2}{\epsilon_{DQD}^-}}}. \quad (7)$$

Where

$$\epsilon_{DQD}^\pm = \pm \epsilon_1 + \sum_{\mathbf{k}} \frac{V_1 V_1^*}{\omega - \epsilon_{\mathbf{k}}} + \frac{\|\pm t_{dots} + \sum_{\mathbf{k}} \frac{V_1 V_2^*}{\omega - \epsilon_{\mathbf{k}}}\|^2}{\omega \pm \epsilon_2 - \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega - \epsilon_{\mathbf{k}}}}, \quad (8)$$

$$T_{\pm} = \pm t_1 \pm t_2 \frac{(\pm t_{dots} + \sum_{\mathbf{k}} \frac{V_1 V_2^*}{\omega - \epsilon_{\mathbf{k}}})}{\omega \pm \epsilon_2 \pm \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega - \epsilon_{\mathbf{k}}}}, \quad (9)$$

and

$$\epsilon_{M2} = \omega - \epsilon_M - \frac{\frac{\omega}{\omega + \epsilon_M} \|t_2\|^2}{\omega - \epsilon_2 - \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega - \epsilon_{\mathbf{k}}}} - \frac{\frac{\omega}{\omega + \epsilon_M} \|t_2\|^2}{\omega + \epsilon_2 - \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega + \epsilon_{\mathbf{k}}}}. \quad (10)$$

On the other hand the spin- \uparrow DOS, which is not coupled to the MZM, can be obtained by taking $t_1, t_2 = 0$, hence giving

$$G_{d_{1\uparrow}, d_{1\uparrow}^\dagger}(\omega) = \frac{1}{\omega - \epsilon_{DQD}^+}. \quad (11)$$

The final results will depend on the broadening parameter of QD i with the lead (Γ_i). This broadening satisfies the equation

$$-i\Gamma_i = \lim_{s \rightarrow 0} \sum_{\mathbf{k}} \frac{V_i^* V_i}{\omega + is - \epsilon_{\mathbf{k}}}. \quad (12)$$

By convention we take Γ_1 as the energy unit for the rest of the project. Finally we compute the DOS

$$\rho_{1\sigma}(\omega) = -\frac{1}{\pi} \text{Im} [G_{d_{1\sigma}, d_{1\sigma}^\dagger}(\omega)]. \quad (13)$$

Similar results can be obtain for the DOS of the second $\rho_{2\sigma}$ by exchanging the indexes 1 and 2 in equation (11).

The density of states provides significant information about the presence of a Majorana zero modes in the dot. The Majorana signature is characterized by a robust zero-mode with height equal to: 1) $\frac{0.5}{\pi\Gamma}$ or 2) The half of the spin- \uparrow DOS. Considering that only the spin- \downarrow channel is connected to the Majorana mode and that the spin- \uparrow and spin- \downarrow channels are decoupled in non-interacting systems, the existence of this second type of Majorana signature is intriguing.

C. Interacting case (NRG)

The Numerical Renormalization Group (NRG)^{12,23,24} is the most successful methods to study interacting quantum impurity models. In this project, the impurity is described by the DQD attached to a Majorana mode. Our code a coulomb repulsion factor of $U = 17.3\Gamma_1$ in both dots and a cut-off energy of $D = 2U = 34.6\Gamma_1$. The spacing with other energy levels is assumed to be higher than D , such that only the two coulomb states are relevant for the system dynamics. Particle-Hole-Symmetry at each dot is obtained when $\epsilon_i = \frac{U}{2}$ in both dots. At this point, each dot has an odd number of electrons, hence, at sufficiently low temperature the system will exhibit

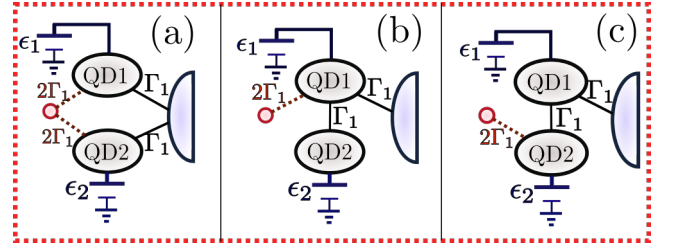


FIG. 2: Cases of study: (a) Symmetric coupling of the DQD to the lead and the MZM. No inter-dot coupling. (b)&(c) Indirect coupling of the second QD through the first dot. The Majorana is coupled to the (b) first dot or to the (c) second dot.

the characteristic Kondo peaks at the Fermi energy Wilson¹². The coexistence of Kondo and Majorana zero modes is still a point of contention in the area and one of the objectives of this part of the project

To improve the efficiency of the code we used the symmetries of the system to maintain a block structure during NRG's iterative diagonalization process. This model preserves the spin- \uparrow particle number \hat{N}_\uparrow and the spin- \downarrow parity $\hat{P}_\downarrow = \pm$ (+ even, - odd). The spin- \downarrow particle number is not preserved due to superconducting-type Majorana coupling (6). The initial Hamiltonian is organized in blocks according to these symmetries. This block structure is preserved during the entire iteration process²⁴. To compute the spectral functions, we use the density matrix renormalization group (DM-NRG)²⁵ in combination with the renown Z-trick method²⁶, which improves spectral resolution at high energies.

III. RESULTS

A. MZM manipulation in non-interacting quantum dots

We call MZM manipulation to the "movements" attributed to the Majorana signature under the tuning of the dot gate voltages (ϵ_1, ϵ_2). This manipulation process is performed in three different set ups that are presented in Fig.2 with definite values of Γ_2 , t_{dots} , t_1 and t_2 . In configuration (a), we coupled the QD symmetrically to the lead and the Majorana mode. With this setup we expect to break the localization of the MZM which should split and tunnel into both dots. In setups (b) and (c) we coupled the second dot indirectly through the first dot. Hence, quantum interference should split the zero mode in two states. Our objective is to observe what occurs with the Majorana signature in this situation. There are two options to connect the MZM in this situation. Attached it directly through the first dot (b) or indirectly through the second dot (c).

The non-interacting results for setups (a),(b) and (c)

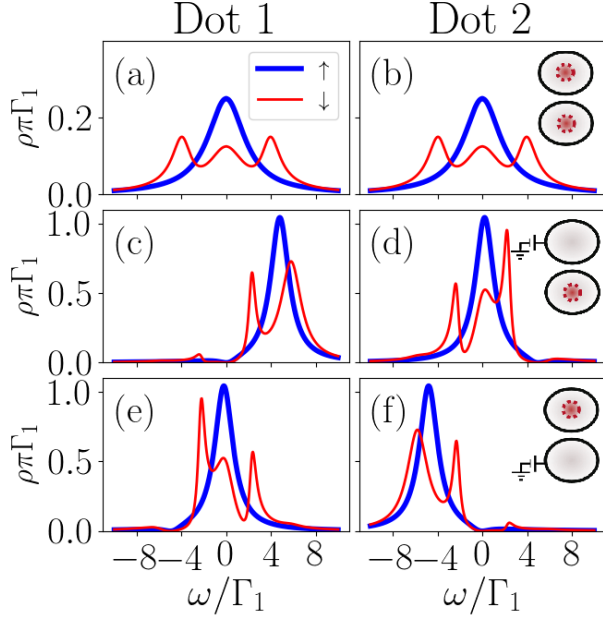


FIG. 3: Non-interaction DOS in the symmetric coupling setup (Fig.2(a)). Bold blue lines: Spin- \uparrow DOS. Thin red lines: Spin- \downarrow DOS.

of FIG.2 are shown at figures FIG.3, FIG.4 and FIG.5 respectively. Each figure depicts the DOS of dot 1(left) and dot 2(right). The gate voltage is initially 0 in both dots at the first row. In the second row the gate voltage is turned on to $\epsilon_1 = 5\Gamma_1$ in the first dot and remains at $\epsilon_2 = 0$ in the second dot. In the third row the first dot's voltage is off $\epsilon_1 = 0$ and we switch on the second dot with a negative voltage of $\epsilon_2 = -5\Gamma_1$. The insets at each row shows which dots exhibit Majorana signatures, depicted by a red dashed circle inside the dot. These images will continuously change under the tuning of gate voltages which represents the manipulation of the Majorana signature.

In FIG.3 we observe the results for the symmetric coupling setup FIG.2(a). In the particle hole symmetric case (first row) the DOS is equal in both dots. Note that that the spin- \downarrow (Thin red line) DOS is the half of the spin- \uparrow (Bold blue line) DOS at the Fermi energy ($\rho_{\downarrow}(0) = \frac{1}{2}\rho_{\uparrow}(0)$). This Majorana signature is similar to the one observed when a single dot is coupled to a Majorana mode.⁸ We may conclude that the Majorana in tunneling inside both dots breaking the localization of the MZM. If a positive or negative gate voltage is induced in one of the dots, as shown in the second and third row of Figure 3(c)-(f), the Majorana zero mode vanishes from that dot. Meanwhile the density of states in the other dot increases while preserving the Majorana signature. This means that the MZM is actually being induced to "leave" this dots and leak into the other dot by the gate voltage activation. This first example of MZM manipulation.

Another example of MZM manipulation occurs when

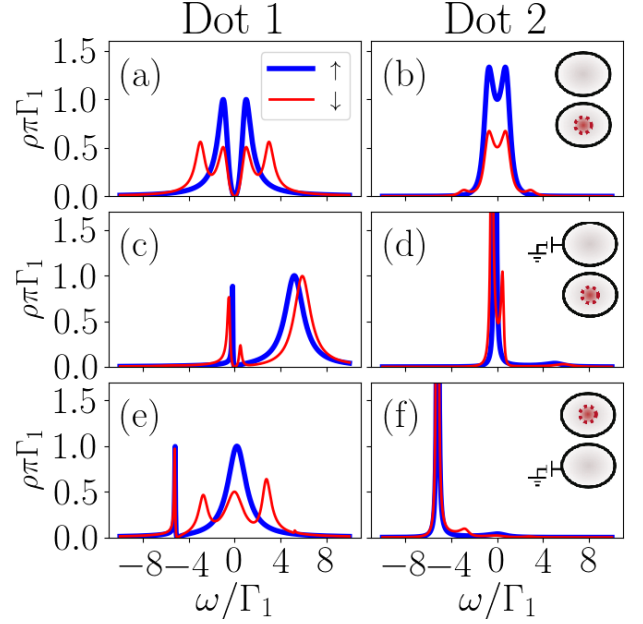


FIG. 4: Non-interaction DOS of the setup in Fig.2(b). Bold blue lines: Spin- \uparrow DOS. Thin red lines: Spin- \downarrow DOS.

the second dot is not directly connected to the lead. In this case, the inter-dot tunneling generates quantum interference which finally destroys the central peak as observe in FIG.4(a) at the spin- \uparrow DOS. The spin- \downarrow channel at FIG.4(a), which is coupled to the MZM, does not exhibit the characteristic Fermi peak either. Instead, the one half Majorana signature at the Fermi energy ($\rho_{\downarrow}(0) = \frac{1}{2}\rho_{\uparrow}(0)$) appears clearly inside the second dot FIG.4(b). This situation prevails when the first dot's gate voltage is turned on FIG.4(c)&(d). While the first dot does not seem to exhibit any type of Majorana signature, the second dot's spin- \downarrow DOS exhibits a robust zero-mode of height $\frac{0.5}{\pi\Gamma}$. The results are more exciting when the second dot's gate voltage is turned on in FIG.4(e)&(f). These figures clearly show how the MZM, previously localized at the second dot, is induced to leave this dot and returned onto the first dot. Moreover, the DOS of spin- \uparrow and spin- \downarrow channels are very similar to the spectral densities observed at FIG.3(d)(e), which means that the previous interference pattern has disappeared due to the presence of this gate voltage.

The results of the third configuration FIG.2(c) appear in FIG.5. Contrary to what was observed in the previous case, this time the Majorana signature is not destroyed by the interference but instead, the $\frac{0.5}{\pi\Gamma}$ -height MZM emerges indirectly in the first dot. This is a perfect way to separate the Majorana's spin- \downarrow DOS from the central spin- \uparrow zero-mode which is still destroyed by the interference. In addition, the second dot still exhibits a Majorana signature as observed in FIG.5(b). In the second row we observe that turning on the gate voltage in dot 1 de-

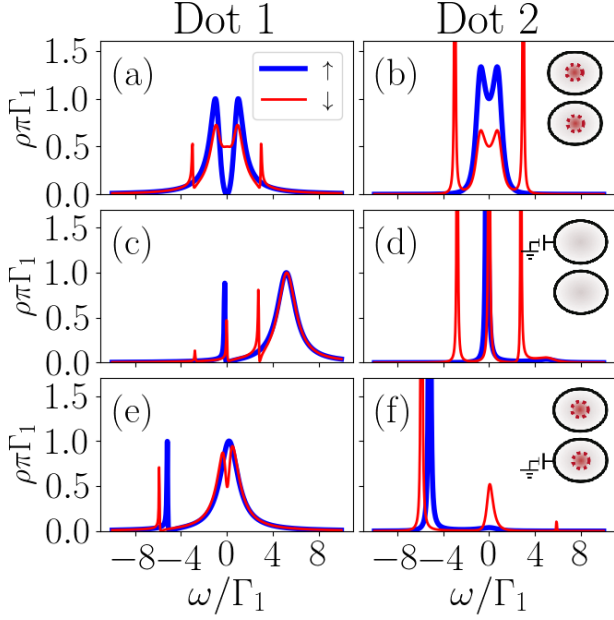


FIG. 5: Non-interaction DOS of the set up in Fig. 2(c). Bold blue lines: Spin-↑ DOS. Thin red lines: Spin-↓ DOS..

stroys the Majorana signature in both dots FIG. 5(c)(d). On the other hand, if the second dot's voltage is switched on the Majorana signature will still appear in both dots, while the spin-↑ quantum interference vanishes.

In total we were able to distinguish two types of Majorana signatures:

- **Type I:** The spin-↓ DOS is the half of the spin-↑ DOS at the Fermi energy ($\rho_{\downarrow}(0) = \rho_{\uparrow}(0)$).
- **Type II:** A spin-↓ zero mode of height $\rho_{\downarrow}(0) = \frac{0.5}{\pi\Gamma_1}$.

B. Majorana manipulation in interacting dots

Now we consider that there is Coulomb repulsion energy of $U = 17\Gamma_1$ in both dots. The factor $\frac{U_i}{2}(\sum_{\sigma} \hat{n}_{i\sigma} - 1)^2$ in (2) favors states with an odd number of electrons (and holes). In addition, particle-hole equilibrium is now achieved when $(\epsilon_{di} + \frac{U_i}{2}) \hat{n}_{i\sigma}$. Any induced gate voltage must be considered as a shifting from this equilibrium point. FIG. 6 shows the DOS at both QDs for the symmetric coupling configuration 2. The two peaks appearing at around $8.6\Gamma_1 = \frac{U}{2}$ represent the two energy levels spaced by the Coulomb repulsion factor U . The central spin-↑ peak is a consequence of the Kondo effect,^{11,12} while the two satellite peaks observed in the inset are the result of the RKKY indirect interaction between both dots.¹⁷⁻¹⁹ Moreover, the system presents a Majorana signature characterized by half spin-↓ DOS at the Fermi energy ($\rho_{\downarrow}(0) = \frac{1}{2}\rho_{\uparrow}(0)$). Note, that in this case the

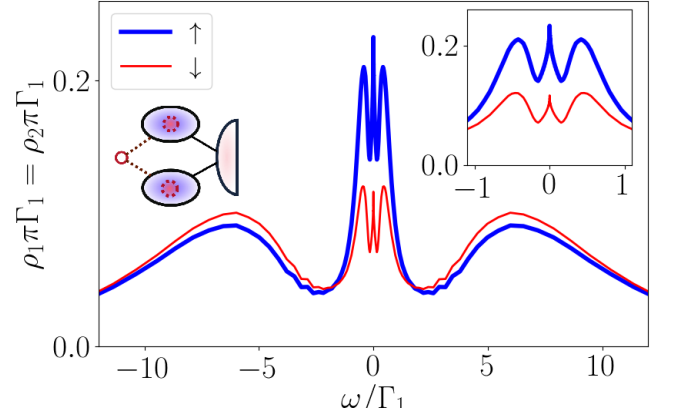


FIG. 6: Density of states of both dots in the symmetric coupling case between the Majorana and the interacting DQD. Bold blue lines: Spin-↑ DOS. Thin red lines: Spin-↓ DOS.

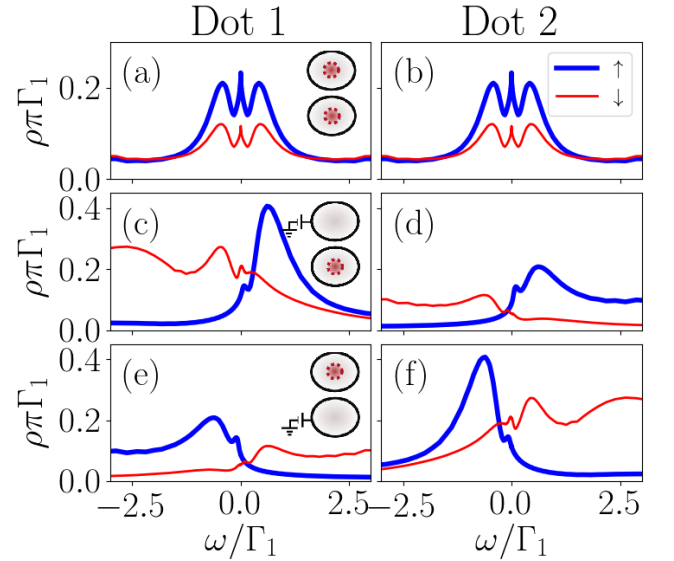


FIG. 7: Density of states in the symmetric coupling arrangement (Fig. 2 first column). Bold blue lines: Spin-↑ DOS. Thin red lines: Spin-↓ DOS.

Majorana signature coexists with the Kondo effect in the DQD as already predicted by Ruiz-Tijerina *et al.* for a single dot.¹⁴

The energy regime we are interested in this project is the one observed in the inset of FIG. 6 close to the Kondo temperature and the emergence of Majorana physics. FIG. 7 shows the NRG results symmetric setup in FIG. 2(a). In agreement with the non-interacting results, we observe that the Majorana signature, which is initially localized in both dots, can be manipulated by tuning the QD gate voltages.

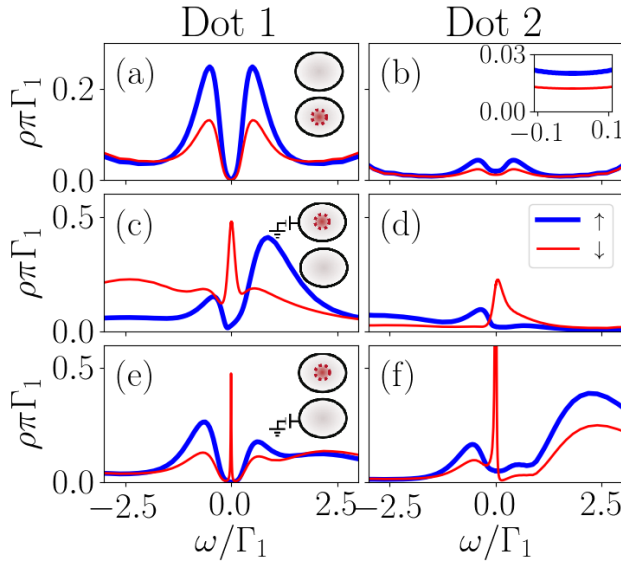


FIG. 8: Density of states in both dots of the case where the only the first QD is attached to both Majorana and Lead (Fig.2 second column) . Bold blue lines: Spin- \uparrow DOS. Thin red lines: Spin- \downarrow DOS.

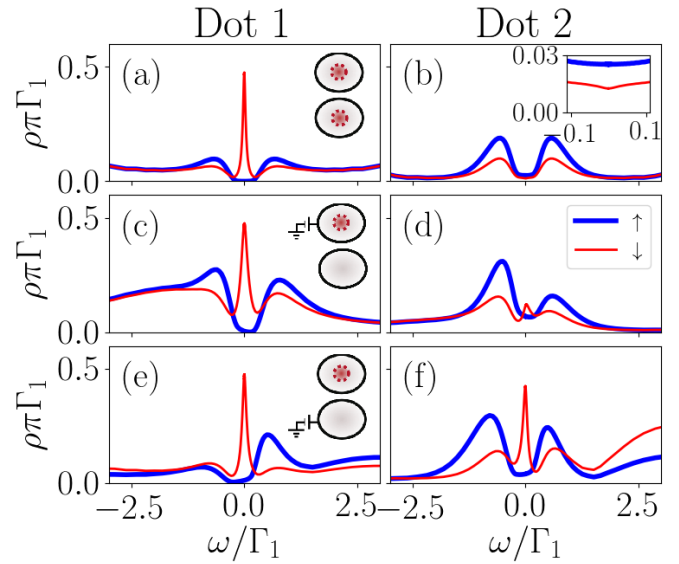


FIG. 9: Density of states of both dots in the case where only the first QD is attached to the lead and the Majorana mode is attached to the second QD. (Fig.2 third column) . Bold blue lines: Spin- \uparrow DOS. Thin red lines: Spin- \downarrow DOS.

IV. CONCLUDING REMARKS

Conclusion goes here.

ACKNOWLEDGMENTS

The authors thank Edson Vernek for enlightening discussions. L.G.G.V.D.S. acknowledges financial support by CNPq (grants No. 307107/2013-2 and 449148/2014-9), and FAPESP (grant No. 2016/18495-4).

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Appendix A: Computation of the Green Function

In Zubarev's fermionic ballistic transport approach²¹ the green function associated to two operators $A(t)$, $B(t)$ is defined as that Fourier transform of the time-ordered anti-commutator of A and B

$$G_{A,B}(\omega) = \mathcal{F} \{ \mathcal{T} [\{ A(t), B(t') \}] \} (\omega). \quad (\text{A1})$$

The Fourier transform of Schrodinger evolution leads to the the transport equations

$$\omega G_{A,B}(\omega) = \delta_{A^\dagger, B} + G_{[A,H],B}(\omega). \quad (\text{A2})$$

Applying this expression to Hamiltonian (1) replacing A and B by the creation and annihilation operators $d_{i\downarrow}^\dagger, f_{i\downarrow}^\dagger, d_i, f, c_k, c_{k\downarrow}^\dagger$ we obtain a linear transport system. To simplify the complexity of the equations we fix $B = d_{1\downarrow}^\dagger$. In addition note that if we replace A by f_{\downarrow} and f_{\downarrow}^\dagger A2 becomes

$$(\omega - \epsilon_M) G_{f_{\downarrow}, d_{1\downarrow}^\dagger}(\omega) = \frac{t}{\sqrt{2}} \left(G_{d_{1\downarrow}, d_{1\downarrow}^\dagger}(\omega) - G_{d_{1\downarrow}^\dagger, d_{1\downarrow}}(\omega) \right) \quad (\text{A3})$$

$$(\omega + \epsilon_M) G_{f_{\downarrow}^\dagger, d_{1\downarrow}^\dagger}(\omega) = \frac{t}{\sqrt{2}} \left(G_{d_{1\downarrow}, d_{1\downarrow}^\dagger}(\omega) - G_{d_{1\downarrow}^\dagger, d_{1\downarrow}}(\omega) \right). \quad (\text{A4})$$

This allows us to take $G_{f_{\downarrow}, d_{1\downarrow}^\dagger}(\omega) = \frac{\omega + \epsilon}{\omega - \epsilon} G_{f_{\downarrow}^\dagger, d_{1\downarrow}^\dagger}(\omega)$. Hence, we can eliminate $G_{f_{\downarrow}, d_{1\downarrow}^\dagger}(\omega)$ from the equations even before we start Gauss-Jordan process.

Writing the other equations we obtain the linear system of the form

$$\mathcal{T} \vec{G}_{d_1^\dagger} = \hat{e}_1, \quad (\text{A5})$$

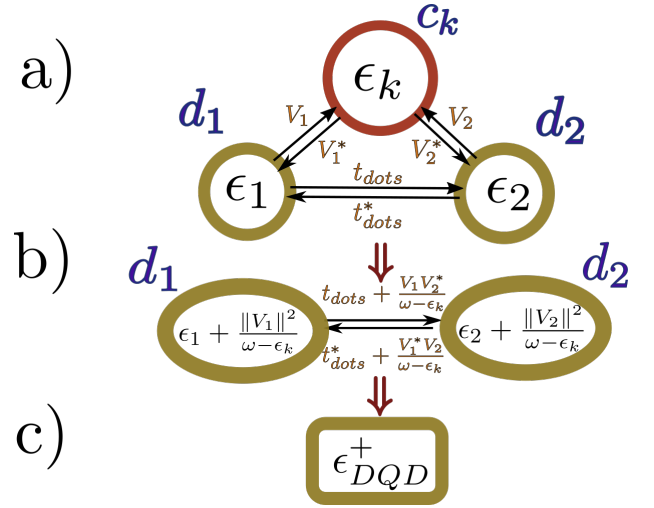


FIG. 10

where \mathcal{T} is the transport matrix

$$\begin{bmatrix} \omega - \epsilon_1 & -V_1^* & -t_{dots} & \frac{-t_1}{\sqrt{2}} & 0 & 0 & 0 \\ -V_1 & \omega - \epsilon_k & -V_2 & 0 & 0 & 0 & 0 \\ -t_{dots}^* & -V_2^* & \omega - \epsilon_2 & \frac{-t_2}{\sqrt{2}} & 0 & 0 & 0 \\ \frac{-\sqrt{2}t_1^*}{\omega + \epsilon_M} & 0 & \frac{-\sqrt{2}t_2^*}{\omega + \epsilon_M} & \omega - \epsilon_M & \frac{\sqrt{2}t_2^*}{\omega + \epsilon_M} & 0 & \frac{\sqrt{2}t_1^*}{\omega + \epsilon_M} \\ 0 & 0 & 0 & \frac{t_2}{\sqrt{2}} & \omega + \epsilon_2 & V_2^* & t_{dots}^* \\ 0 & 0 & 0 & 0 & V_2 & \omega + \epsilon_k & V_1 \\ 0 & 0 & 0 & \frac{t_1}{\sqrt{2}} & t_{dots} & V_1^* & \omega + \epsilon_1 \end{bmatrix}, \quad (\text{A6})$$

$\vec{G}_{d_1^\dagger}$ is the column vector

$$[G_{d_{1\downarrow}, d_{1\downarrow}^\dagger}(\omega), G_{c_{k\downarrow}, d_{1\downarrow}^\dagger}(\omega), G_{d_{2\downarrow}, d_{1\downarrow}^\dagger}(\omega), G_{f_{\downarrow}, d_{1\downarrow}^\dagger}(\omega), G_{d_{2\downarrow}^\dagger, d_{1\downarrow}^\dagger}(\omega), G_{c_{k\downarrow}^\dagger, d_{1\downarrow}^\dagger}(\omega), G_{d_{1\downarrow}^\dagger, d_{1\downarrow}^\dagger}(\omega)]^T$$

and \hat{e}_1 is the vector with entries $\hat{e}_{1_n} = \delta_{1_n}$.

The graph associated to this matrix is the one in FIG.11. The energies inside each vertex are given by subtracting the corresponding diagonal term from ω . The couplings are just the negative of the off-diagonal terms.

1. The double quantum dot

To explain the process of Gaussian elimination we will obtain the green function for the case without Majorana fermion ($t_1 = t_2 = 0$). The transport matrix for this system is

$$\begin{bmatrix} \omega - \epsilon_1 & -V_1 & -t_{dots} \\ -V_1^* & \omega - \epsilon_k & -V_2 \\ -t_{dots}^* & -V_2^* & \omega - \epsilon_2 \end{bmatrix}. \quad (\text{A7})$$

The graph associated to this matrix can be observed in FIG10.a). To eliminate the vertex c_k we just need to

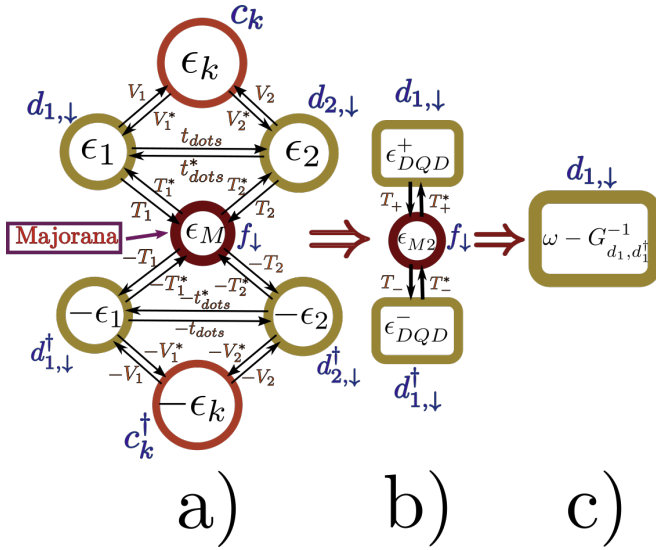


FIG. 11: Transport flow in a DQD Majorana system.

subtract from (A7) the rank-1 matrix that cancels the row and the column corresponding to c_k . This matrix is

$$\begin{bmatrix} \frac{V_1^* V_1}{\omega - \epsilon_k} & -V_1^* & \frac{V_2 V_1^*}{\omega - \epsilon_k} \\ -V_1 & \omega - \epsilon_k & -V_2 \\ \frac{V_2^* V_1}{\omega - \epsilon_k} & -V_2^* & \frac{V_2^* V_2}{\omega - \epsilon_k} \end{bmatrix}. \quad (\text{A8})$$

The result of (A7) - (A8) is

$$\begin{bmatrix} \omega - \epsilon_1 - \frac{V_1^* V_1}{\omega - \epsilon_k} & 0 & -t_{dots} - \frac{V_2 V_1^*}{\omega - \epsilon_k} \\ 0 & 0 & 0 \\ -t_{dots}^* - \frac{V_2^* V_1}{\omega - \epsilon_k} & 0 & \omega - \epsilon_2 - \frac{V_2 V_1^*}{\omega - \epsilon_k} \end{bmatrix} \quad (\text{A9})$$

which is depicted by the graphs in FIG.10.b). The next step is to pop-out the vertex d_2 following the same procedure. At the end, the energy inside the vertex d_1 will be

$$\epsilon_{DQD}^+ = \epsilon_1 + \sum_{\mathbf{k}} \frac{V_1 V_1^*}{\omega - \epsilon_{\mathbf{k}}} + \frac{\|t_{dots} + \sum_{\mathbf{k}} \frac{V_1 V_2^*}{\omega - \epsilon_{\mathbf{k}}}\|^2}{\omega - \epsilon_2 - \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega - \epsilon_{\mathbf{k}}}} \quad (\text{A10})$$

and the green function of $G_{d_1 d_1^\dagger}(\omega)$ in a DQD will be given by $\frac{1}{\omega - \epsilon_{DQD}}$ (see FIG.10.c)).

2. Solution of the transport equations

The previous procedure can be generalized into the following algorithm:

1. Computing the transport equations with the second term fixed in the creation operator of the dot.
2. Setting up the graph associated to the transport system.

3. Popping out the vertexes of the graph. Each popping process carries the following steps.

- (a) Computing the extra-terms in the energies and couplings based on the walks passing through the popped vertex.
- (b) Eliminating this vertex from the graph.
- (c) Iterating till there is only one vertex.

4. The energy in the remaining vertex d is $\epsilon_d = \frac{1}{\omega - G_{d, d^\dagger}(\omega)}$.

Following these steps it is possible to solve the general case. We start with the graph in FIG.11 and we pop out the vertexes $c_k, c_k^\dagger, d_{2,\downarrow}$ and $d_{2,\downarrow}^\dagger$ in that order. The energies associated to $d_{1,\downarrow}$ and $d_{1,\downarrow}^\dagger$ will be similar to (A10) giving

$$\epsilon_{DQD}^\pm = \pm \epsilon_1 + \sum_{\mathbf{k}} \frac{V_1 V_1^*}{\omega - \epsilon_{\mathbf{k}}} + \frac{\|\pm t_{dots} + \sum_{\mathbf{k}} \frac{V_1 V_2^*}{\omega - \epsilon_{\mathbf{k}}}\|^2}{\omega \pm \epsilon_2 - \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega - \epsilon_{\mathbf{k}}}}. \quad (\text{A11})$$

There is also a correction in the couplings between the Majorana mode and $d_{1,\downarrow}, d_{1,\downarrow}^\dagger$ given by

$$T_\pm = \pm t_1 \pm t_2 \frac{(\pm t_{dots} + \sum_{\mathbf{k}} \frac{V_1 V_2^*}{\omega - \epsilon_{\mathbf{k}}})}{\omega \pm \epsilon_2 \pm \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega - \epsilon_{\mathbf{k}}}}. \quad (\text{A12})$$

Finally since the Majorana is in contact with dot 2, there is an extra-term appearing in the Majorana energy given by

$$\epsilon_{M2} = \omega - \epsilon_M - \frac{\frac{\omega}{\omega + \epsilon_M} \|t_2\|^2}{\omega - \epsilon_2 - \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega - \epsilon_{\mathbf{k}}}} - \frac{\frac{\omega}{\omega + \epsilon_M} \|t_2\|^2}{\omega + \epsilon_2 - \sum_{\mathbf{k}} \frac{V_2 V_2^*}{\omega + \epsilon_{\mathbf{k}}}}. \quad (\text{A13})$$

With all the terms of the graph in FIG.11.b) computed, it only remains to pop out vertexes d_1^\dagger and f_\downarrow in that order to obtain the result in equation (11).

$$G_{d_{1\downarrow}, d_{1\downarrow}^\dagger}(\omega) = \frac{1}{\omega - \epsilon_{DQD}^+ - \frac{\|T_+\|^2}{\omega - \epsilon_{M2} - \frac{\|T_-\|^2}{\epsilon_{DQD}^-}}}. \quad (\text{A14})$$