

Negative hybridization: A potential cure for braiding with imperfect Majorana modes

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Majorana zero modes, the elementary building blocks for the quantum bits of topological quantum computers, are known to suffer from hybridization as their wavefunctions begin to overlap. This breaks the ground-state degeneracy, splitting their energy levels and leading to an accumulation of error when performing topological quantum gates. Here, we show that the energy splitting of the Majorana zero modes can become *negative*, which can be utilized to reduce the average hybridization energy of the total gate. We present two illustrative examples where *negative hybridization* suppresses gate errors to such an extent that they remain below the fault-tolerance threshold. As an intrinsic property of Majorana zero modes, negative hybridization enables systems based on imperfect Majorana zero modes to regain functionality for quantum information processing.

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Introduction. Topological quantum computing promises fault tolerance against local perturbations, as quantum gates are enacted by pairwise exchanges of spatially separated Majorana modes, referred to as braids [1–7]. A topological qubit consists of four Majorana zero modes (MZMs), and the nonlocality of such a qubit is key for its topological protection [8,9]. A major research question revolves around the actual fault tolerance of a Majorana-based quantum computer [10]. It is well established that the successful execution of a quantum gate necessitates that the corresponding braiding time, T , satisfies the “speed limit” of Majorana modes [3,11,12]:

$$\frac{\hbar}{\Delta_{\text{topo}}} \ll T \ll \frac{\hbar}{E_{\text{hyb}}}, \quad (1)$$

where Δ_{topo} is the size of the topological gap and E_{hyb} is the hybridization energy of a pair of Majorana modes.

The lower speed limit, $\hbar/\Delta_{\text{topo}}$, protects the Majorana subspace from transitions into, or coupling to, bulk states. Intuitively, if the Majorana modes are moved too fast, they may undergo transitions to higher (excited) states, thus losing quantum information. The smaller the gap, the greater the likelihood of this diabatic error occurring [13–15]. The obvious way to escape from such errors is to perform the dynamics of the Majorana modes sufficiently slowly, i.e., adiabatically.

The source of the upper speed limit, $\hbar/E_{\text{hyb}}$, is less obvious; the hybridization between two Majorana modes, stemming from the spatial overlap of their wavefunctions, causes a splitting of their energies away from zero. The accumulated hybridization over the braid time leads to rotations on the Bloch sphere of the topological qubit: so-called Majorana oscillations [16,17]. In other words, the larger the braid time, the larger the magnitude of qubit errors. This effect

also occurs in static systems where the MZMs hybridize too strongly, removing their topological protection and non-Abelian statistics [17]. Although Majorana modes are able to reach exactly zero energy for finely tuned systems [18], only in the thermodynamic limit can they exist at zero energy in the presence of perturbations. Thus, any finite system will eventually be confronted with the upper speed limit and Majorana oscillations.

The conundrum of a topological quantum computer can hence be formulated as finding a way to satisfy the lower and upper speed limits of Majorana modes simultaneously, as given in Eq. (1). This will be difficult to achieve for many of the current experimental systems that claim to host Majorana modes, due to material and device constraints such as limited lengths and small superconducting gaps.

In this Letter, we introduce the concept of *negative hybridization*, a fundamental property of Majorana modes, which can be thought of as an intrinsic error correction property of Majorana modes. We argue that negative hybridization can be used to drastically reduce E_{hyb} (or even set $E_{\text{hyb}} \rightarrow 0$) and thus push the upper speed limit $\hbar/E_{\text{hyb}}$ to larger times (or even send it to infinity). In other words, it is possible for MZMs “without topological protection” to have their non-Abelian statistics restored. This is emphasized when we introduce a *symmetry-enforced braiding protocol* that utilizes negative hybridization to produce corrected braids, which are robust against perturbations, vital for stable quantum information processing.

Negative hybridization. In the following, we consider a topological superconductor hosting a Majorana bound state (MBS), i.e., a zero-energy excitation consisting of two MZMs that are spatially separated. We can identify this excitation by analyzing the Bogoliubov–de Gennes (BdG) Hamiltonian in its diagonal representation, given by $H = \sum_i E_i d_i^\dagger d_i$. Here, the operators $d_i^\dagger$ correspond to the creation of a Bogoliubov quasiparticle, which will increase the energy of the system by $E_i \geq 0$. For the infinite system, the lowest-energy quasiparticle, $d_1^\dagger$, has energy $E_1 = 0$. Furthermore, we can use

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these associated quasiparticle operators to construct Majorana operators:

$$\gamma_1 = d_1^\dagger + d_1, \quad \gamma_2 = i(d_1^\dagger - d_1). \quad (2)$$

Each of these operators is its own particle-hole conjugate, corresponding to an MZM exponentially localized at a boundary of the topological region. Together, all γ_i satisfy a Clifford algebra. Exchanging two MZMs in a clockwise fashion acts on the low-energy manifold, consisting of the ground state and the occupied MBS state, with the braiding operator

$$B_{12} = B_{21}^{-1} = \exp\left(\frac{\pi}{4}\gamma_1\gamma_2\right). \quad (3)$$

Critically, these braids allow for quantum Clifford gates to be enacted purely through the exchange of MZMs, forming the basis for *topological quantum computation*.

Since Majorana modes in finite systems are never truly at zero energy, it is common to quantify the effect of energy splitting of the Majorana modes on the performance of the resulting gate by considering the fidelity

$$\mathcal{F} = |\langle \Psi_{\text{target}} | \Psi_{\text{final}} \rangle|^2, \quad (4)$$

where $|\Psi_{\text{final}}\rangle = U(T)|\Psi_{\text{initial}}\rangle$ and the implemented braiding protocol is encoded in the time-evolution operator $U(T)$. By this measure, a perfect braid would have a fidelity of 1. We assert that achieving a fidelity of $\geq 99\%$, exceeding the threshold required for error correction [19,20], over an extended regime of total braid times T (i.e., no requirement of *temporal* fine-tuning) is sufficient to demonstrate non-Abelian statistics.

When MZMs are kept at a constant distance throughout a braid, possible for MZMs pinned to vortices in a topological superfluid or superconductor [2], their energy splitting over time can remain constant [17]. However, when considering other platforms such as wire networks where MZMs are physically moved (e.g., by means of gate protocols), the separation, and thus energy splitting, between MZMs will change over time [12,17,21,22]. Taking the simplest example, we consider two MZMs on a T-junction geometry performing a $\sqrt{Z}$ gate [23] in which the two MZMs change their position once, as illustrated in Fig. 1(a). As expected, when MZMs get closer to each other, the energy splitting increases [see Fig. 1(b) and Movie S1 in the Supplemental Material as examples [24]]. The relevant quantity describing this splitting over time is the average hybridization energy:

$$\bar{E} = \frac{1}{T} \int_0^T E_1(t) dt. \quad (5)$$

For the simplest case of two MZMs, this hybridization causes additional rotations between the MZM states, on top of the expected braid statistics, and is given by

$$U_{\text{hyb}}(T) = \exp\left(-\frac{\bar{E}T}{2}\gamma_1\gamma_2\right). \quad (6)$$

This factor causes the resulting fidelity of a $\sqrt{Z}$ gate to oscillate as $\mathcal{F}(T) = \cos^2(\bar{E}T/2)$ [12,21]; the smaller $\bar{E}$ is, the longer the fidelity remains at, or close to, 1.

The spectra shown in Figs. 1(b), 1(d), and 1(f) are obtained by diagonalizing the BdG Hamiltonian where, by

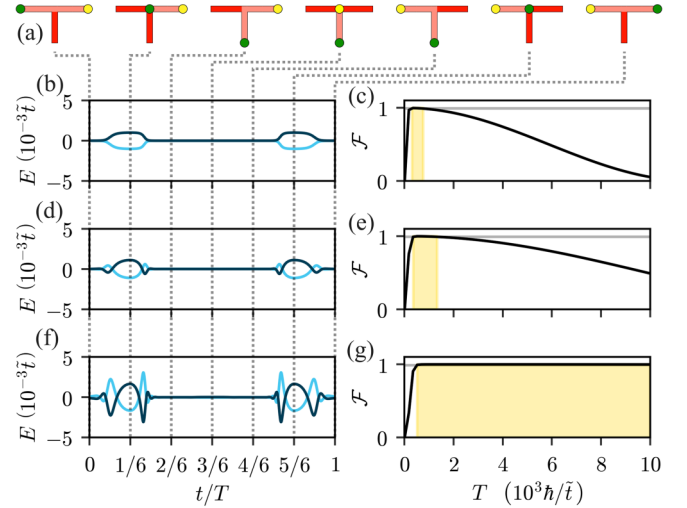


FIG. 1. Energy hybridization of two MZMs performing a $\sqrt{Z}$ gate. (a) Sketch of two MZMs (colored dots) being exchanged on a T-junction geometry, performing a $\sqrt{Z}$ gate. Light and dark red sections correspond to topological and trivial parameters, respectively. (b) Example of energy splitting between the two MZMs as a function of time t (with total braid time T). (c) Fidelity $\mathcal{F}$ corresponding to panel (b) for different braid times T . The braid times where the fidelity exceeds the threshold of 99% to be error correctable are highlighted in yellow. (d) and (e) Same as panels (b) and (c) but with slightly different parameters, leading to weak negative hybridization. $\mathcal{F}$ vs T shows enlarged yellow region. (f) and (g) Same as panels (b) and (c) and panels (d) and (e) but with parameters that induce significant negative hybridization and thus vanishing $\bar{E}$. As a consequence, Majorana oscillations are absent and $\mathcal{F}$ exceeds 99% irrespective of T .

construction, the doubling of the degrees of freedom leads to identical positive and negative energies. Typically, one would only consider the positive half of such spectra as physical. Yet, when inspecting spectra similar to the examples in Figs. 1(d) and 1(f), it is clear that the energy of the MBS crosses through $E = 0$ and spans both positive and negative values.

This energy crossing corresponds to a swapping of the ground-state parity, where the original many-body ground state becomes higher in energy than the first excited state. Since the two states differ by one quasiparticle, the parity must change (see the Supplemental Material [24]). These parity swaps can be calculated by taking the sign of the Pfaffian of the BdG Hamiltonian in the Majorana basis [25]. This confirms the zero-energy crossings in Figs. 1(d) and 1(f).

As a consequence of this effect, when integrating the energy splitting to calculate $\bar{E}$, hybridization of the positive sections will cancel with the negative sections, thereby reducing the magnitude of $\bar{E}$. This can be seen in Figs. 1(b), 1(d), and 1(f) where their corresponding fidelities [Figs. 1(c), 1(e), and 1(g)] exceed the fault-tolerance threshold for progressively wider ranges of braid times due to the reduced hybridization error. We refer to the phenomenon of $E_1(t)$ changing sign, leading to an improvement in braid fidelity, as *negative hybridization*. Due to the presence of negative hybridization, the requirement that a “good” braid has $\bar{E} = 0$ does not imply that E_1 must remain at zero for the entire

process, suggesting that accurate gates may be achievable even in systems where hybridization is unavoidable.

The remainder of this Letter will explore ways in which the braid protocol may be adapted to induce negative hybridization, thus improving quantum gates afflicted with fidelities far below the required accuracy ($<99\%$) to far above ($>99\%$), thus restoring fault tolerance [19,20].

Negative hybridization enforced by a local gate. In the following, we will investigate a simple scenario for the utilization of negative hybridization, which is potentially controllable and implementable in various platforms. To this end, we study the simplest topological superconductor model, the Kitaev chain. All of the simulations were performed on Kitaev chains that are combined into multi-T-junction geometries, consisting of multiple vertical legs connected to a horizontal chain and can be described by

$$\begin{aligned}
 H(t) = & \sum_{i=0}^{2M-1} \left[\sum_{j=1}^{L_i} -\mu_{i,j}(t) c_{i,j}^\dagger c_{i,j} \right. \\
 & + \sum_{j=1}^{L_i-1} (-\tilde{t} c_{i,j}^\dagger c_{i,j+1} + \Delta e^{i\phi_i} c_{i,j}^\dagger c_{i,j+1}^\dagger + \text{H.c.}) \Big] \\
 & + \sum_{i=1}^{2M-1} (-\tilde{t} c_{0,x_i}^\dagger c_{i,1} + \Delta e^{i\phi_i} c_{0,x_i}^\dagger c_{i,1}^\dagger + \text{H.c.}). \quad (7)
 \end{aligned}$$

Here, M is the number of MZM pairs and L_i is the number of sites in the i th leg, ordered such that the horizontal backbone is L_0 and the vertical legs, from left to right, are L_1 through L_{2M-1} . As the MZMs are evenly spaced along the horizontal chain, it is convenient to use half that distance, L_h , and the number of buffer sites, L_{buffer} , to express the horizontal length: $L_0 = 2L_{\text{buffer}} + (2M-1)(2L_h+1)$. Each pair of adjacent MZMs has a vertical leg between them so that they can be braided without the movement of any of the other MZMs. These are attached to the horizontal leg at sites $x_i = L_{\text{buffer}} + (2i-1)L_h + i$. To ensure the swaps in parity that the symmetric gates are built around, the lengths of the vertical legs alternate such that for $n \in \mathbb{Z}^+$, $L_{2n-1} = L_v$ and $L_{2n} = L_v + 1$, where L_v is the length of the odd numbered legs. To mirror the $p_x + ip_y$ superconductivity that provides the correct braiding statistics, we define the superconducting phases of each section such that $\phi_0 = 0$ and $\phi_{i>0} = \frac{-\pi}{2}$ so that all of the vertical legs ($i > 0$) are orthogonal to the horizontal backbone ($i = 0$). Further detail on the system setup, including the use of buffer sites, is given in the Supplemental Material [24].

MZMs can be moved by carefully changing the chemical potentials of certain sites between a topological ($|\mu/\tilde{t}| < 2$) and a trivial ($|\mu/\tilde{t}| > 2$) value, which we refer to as μ_{topo} and μ_{triv} , which take the values $\mu_{\text{topo}} = 0.5\tilde{t}$ and $\mu_{\text{triv}} = 5\tilde{t}$ for all simulations unless otherwise specified. For details about the ramping protocol used, see the Supplemental Material [24].

We first consider a T-junction geometry where an additional local gate potential with strength V , satisfying $\mu_{\text{topo}} < V < \mu_{\text{triv}}$, is applied to the top site of the vertical leg [see inset of Fig. 2(b)]. For simplicity, we restrict this gate to a single site (for extensions to more than a single site [23,26],

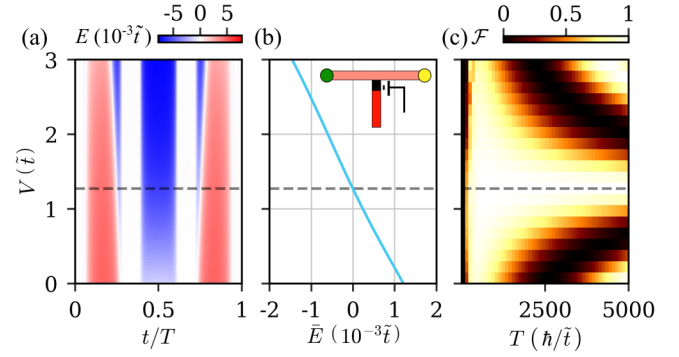


FIG. 2. Negative hybridization induced by a local gate potential. ($M, L_h, L_v, L_{\text{buffer}}, \Delta$) = (1, 5, 11, 0, $0.7\tilde{t}$). (a) The hybridization energy between the MZMs throughout the braid, for varying applied voltage strengths, V , located at the site below the junction. As V increases, the main parity swaps occur earlier in the braid, leading to greater amounts of negative hybridization. (b) Integrated energy, $\bar{E}$, for the braid as a function of V . The dashed line at $V = 1.27\tilde{t}$ marks the potential where the braid has $\bar{E} = 0$. Inset: T-junction with the additional local gate highlighted in black. (c) Fidelity of the braid for various V and T . Around $V = 1.27\tilde{t}$, the fidelity remains close to 1, even for large values of T .

see the Supplemental Material [24]), modifying μ locally, i.e., $\mu_{\text{site}}(t) \rightarrow \mu_{\text{site}}(t) - V$.

At the start of the braid, the site has chemical potential $\mu_{\text{site}}(0) = \mu_{\text{triv}} - V > 0$, but during the braid there are times when the bottom leg becomes topological and $\mu_{\text{site}}(t) = \mu_{\text{topo}} - V < 0$, changing sign. Due to the chosen location of the potential, this change in sign occurs when the MZMs are strongly hybridizing. In the limit $|\Delta| = |\tilde{t}|$, the parity simply becomes the product of the chemical potentials on all sites of the chain. Despite the number and timing of parity swaps changing as the magnitude of Δ is lowered, it remains a valid approximation, as shown in the Supplemental Material [24]. Therefore, if both μ_{topo} and μ_{triv} have the same sign, the parity of the system remains constant throughout a braid. However, by changing the sign of the chemical potential on just one site, we are able to induce negative hybridization.

As seen in Fig. 2(a), where the hybridization energy throughout a $\sqrt{Z}$ gate is plotted as a function of time and gating potential, increasing V leads to an earlier parity swap as μ_{site} is brought closer to zero, thus introducing more negative hybridization. From the hybridization energy of each braid, we obtain $\bar{E}(V)$, shown in Fig. 2(b). As expected, increasing the amount of negative hybridization by increasing V leads to a reduction of $\bar{E}$, with $\bar{E} = 0$ for $V \approx 1.27\tilde{t}$. In Fig. 2(c) the fidelity of the braid is plotted for varying braid times, T . These show the expected Majorana oscillations for the clean system, limiting the range of T for successful braids, but as V approaches $1.27\tilde{t}$, the region where $\mathcal{F} > 0.99$ rapidly extends to large times. At this voltage, the fidelity is essentially independent of braid time, leaving only the contributions from the braiding statistics, thus raising the gate fidelity well above the fault-tolerance threshold.

It should be noted that the strength of the potential itself modifies the magnitude of MZM hybridization [27]. Therefore, the value of V that minimizes $\bar{E}$ is difficult to determine a

priori. However, calibration through repeating the same braid for changing values of V seems to be a rather simple task, and as demonstrated by Fig. 2, even an imperfectly calibrated gate potential can significantly enhance the braiding performance. We stress that for the quantum dot platform [18,28,29] that hosts “poor man’s Majoranas,” our described protocol for utilizing negative hybridization could be directly implemented.

Negative hybridization due to symmetric braids. The above reduction of $\bar{E}$ requires the calibration of a system parameter in order to significantly reduce the dynamic phase accumulated over the braiding process, with perfect phase cancellation necessitating fine-tuning. Thus, for our second example, we use the original braid as the basis of the new protocol, where the second half of the braid is built to have the same magnitude of hybridization but opposite sign, leading to $\bar{E} = 0$. The key is to utilize the symmetry properties of topological superconductors, here specifically the Kitaev model. Similar reasoning has been proposed, where X gates are acted on ancillary modes to induce parity swaps [30,31]; however, our method relies on the intrinsic parity switching of the MZMs and thus does not rely on external gates or zero modes.

The antisymmetry of the spectrum of the Kitaev chain for odd chain lengths [32,33] with respect to positive and negative chemical potentials allows us to perform identical operations but with opposite parity. This method can be extended to other spinful topological superconductor models [34–36] as shown in the Supplemental Material [24]. We stress that restricting segments of a T-junction to an odd number of lattice sites or atoms is unproblematic for quantum dot arrays [28,29] and for magnet-superconductor hybrid structures, where atomic chains are built atom by atom [37–39].

For the proposed method to work, two challenges need to be overcome: (1) The additional operations must not contribute to the braiding statistics. (2) The process that induces the parity swap *cannot* create uncanceled hybridization itself.

To produce a trivial movement of the MZMs that accumulates the same hybridization as a braid, yet acts as an identity gate, we first decompose the original braid into two halves: $B_{ij}^{(1)}$ and $B_{ij}^{(2)}$. For illustration purposes, suppose that the first operation corresponds to the process $t = 0$ to $t = 3/6T$, depicted in Fig. 1(a), while the second operation corresponds to the process $t = 3/6T$ to $t = T$. The original braid is therefore defined as $B_{ij}^{(2)}B_{ij}^{(1)}$, while the new trivial operations are given by either $B_{ij}^{(1)}B_{ij}^{(1)-1}$ or $B_{ij}^{(2)}B_{ij}^{(2)-1}$.

Figure 3(a) depicts the state of the system as a corrected $\sqrt{X}$ gate is performed, where the dots represent MZMs, and the color of each section denotes the chemical potential as given by the color bar. The time-dependent hybridization energy is plotted in Fig. 3(b), where each sixth of the braid is colored and labeled according to the process occurring. At the midpoint of the uncorrected braid, $T = 1/6$, all sites are topological with $\mu = \mu_{\text{topo}}$ and the entire system can be simultaneously ramped from μ_{topo} to $-\mu_{\text{topo}}$, defining the operation $M_{ij}^{(+)}$ whose counterpart, $M_{ij}^{(-)}$, transforms the system back from a negative chemical potential to a positive one. As the $M_{ij}^{(\pm)}$ protocols are symmetric in μ , they have no net effect on $\bar{E}$, as seen in Fig. 3(b).

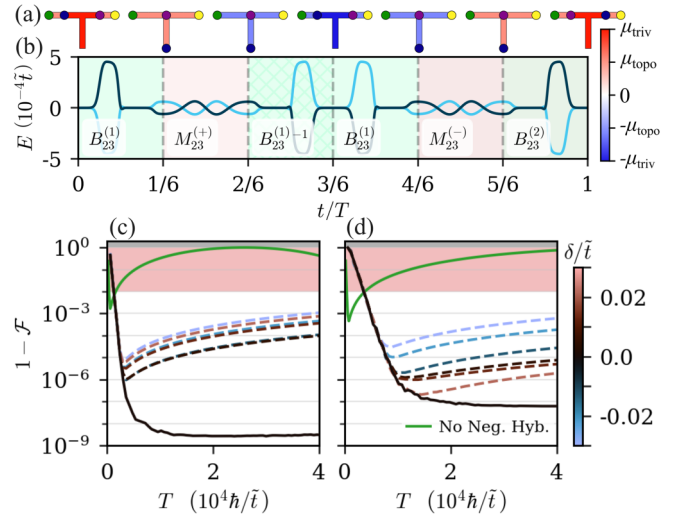


FIG. 3. Symmetric braids: errors of corrected and uncorrected single-qubit gates. $(M, L_h, L_v, L_{\text{buffer}}, \Delta) = (2 \text{ and } 1, 8, 8, 0, 0.8\bar{t})$. (a) Schematic of the corrected $\sqrt{X}$ gate showing MZMs (colored) at the borders of the topological (light red and light blue) and trivial (dark red and dark blue) regions. The red regions correspond to positive chemical potentials, while the blue ones indicate that the chemical potential is negative. (b) Energy spectrum of the Majorana bound states for the corrected $\sqrt{X}$ gate as a function of time, labeled with the braid operations taking place. (c) Error in fidelity for $\sqrt{X}$ gate vs braid times. The red region covers the fidelities that exceed the 1% error threshold. The green curve is the uncorrected result; all other lines have been corrected by means of symmetric braids. $\delta = 0$ (solid black) is perfectly symmetric, and finite δ values (dashed) indicate deviation from perfect symmetry. (d) Same as panel (c) but for the $\sqrt{Z}$ gate.

We note that swapping the chemical potentials for a braid in a $\sqrt{Z}$ gate is more difficult as there are always regions of the junction that are trivial. These cannot directly go from μ_{triv} to $-\mu_{\text{triv}}$ without causing the formation of new MZMs—effectively introducing quasiparticle poisoning. To avoid this, we employ a more sophisticated procedure, as outlined in the Supplemental Material [24].

By taking advantage of the parity’s antisymmetric structure with respect to the chemical potential, the routines for the corrected braids are agnostic to the specific parameters of the system; thus, we do not need any prior knowledge of the hybridization dynamics to produce a braid with no net hybridization.

To demonstrate the effectiveness of these protocols, we simulate the corrected braids corresponding to $\sqrt{X}$ and $\sqrt{Z}$ gates. The resulting error in the fidelity, $1 - \mathcal{F}$, as a function of total braiding time, T , is shown in Figs. 3(c) and 3(d), compared to the error of the conventional, *uncorrected* gates. The red region shows errors larger than 1%, where quantum computing fails. The uncorrected (solid green) curve briefly escapes the red region, but then returns due to the Majorana oscillations. Meanwhile, the corrected (solid black) curve decreases monotonically with the diabatic error before plateauing at an exceptionally small error $1 - \mathcal{F} < 10^{-7}$, far below the level required for error correction.

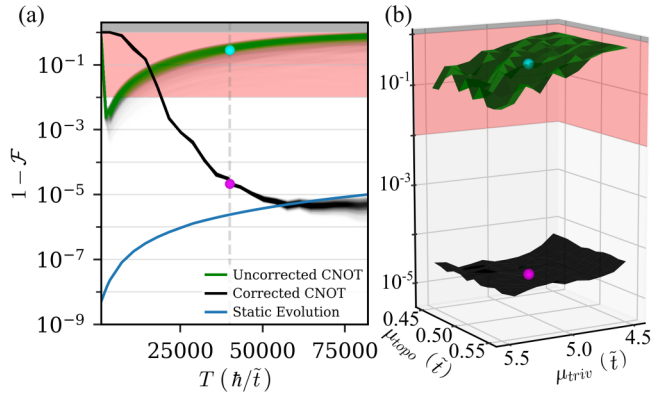


FIG. 4. CNOT gate error. $(M, L_h, L_v, L_{\text{buffer}}, \Delta) = (3, 8, 8, 4, 0.8\tilde{t})$. (a) Loss of fidelity for the corrected and uncorrected CNOT gates vs braid time, compared to the natural time evolution of the system. Results slightly depend on initial states, leading to the spread of the uncorrected curve. (b) Average fidelities for the corrected and uncorrected gates over a range of chemical potential values, demonstrating the robustness of the protocol. Results are for the braid time $T = 4000\hbar/\tilde{t}$, indicated by the dashed line in panel (a). The dots indicate common data between panels (a) and (b).

As it is unrealistic to assume arbitrarily precise control over the braid, we introduce the parameter $\delta = \mu_{\text{topo}}^+ + \mu_{\text{topo}}^-$, which quantifies the asymmetry between the positive and negative chemical potentials. The color of the dashed lines in Figs. 3(c) and 3(d) corresponds to specific values of δ . In general, we see that although this imperfection reintroduces hybridization dependent error, it remains orders of magnitude lower than the uncorrected case, even when δ is a considerable fraction of μ_{topo} . This highlights the robustness of the corrected gates, showing they are able to function without fine-tuning.

Individual corrected gates can be easily combined and extended beyond a single qubit. To this end, we demonstrate the CNOT gate on a two-qubit system, comprised of six MZMs. We note that the further extension to more than two qubits is straightforward, e.g., via projective measurements [40,41], where it is sufficient to correct single- and two-qubit gates.

The CNOT gate in the even-parity subspace can be expressed as $B_{12}B_{23}B_{12}B_{65}B_{34}B_{21}B_{12}B_{23}B_{12}$ [42,43]. Therefore, the corrected CNOT gate is simply a collection of multiple braids, each being individually corrected. In Fig. 4(a), we show the braiding error, $1 - \mathcal{F}$, of the CNOT gate. As the

fidelity of the CNOT gate depends on the initial state, we test multiple random initial states, shown green for the uncorrected and black for the corrected gates. The improvement due to negative hybridization is striking, showing that the proposed method to produce corrected braids can indeed be extended and applied to complex circuits with multiple qubits. The blue line in Fig. 4(a) is the error of the static system, i.e., the accumulated error from the hybridization of MZMs simply waiting in their initial positions. Although much smaller than the error during the braids where the MZMs move closer to each other, it is neglected in the correction process and thus forms a lower bound on the error of the corrected gates. In Fig. 4(b), we show deviations from the parameters μ_{topo} and μ_{triv} used in Fig. 4(a) at time $T = 4000\hbar/\tilde{t}$, indicated by the dots and dashed line. This demonstrates the robustness of the protocol with respect to different choices of parameters. The errors of corrected and uncorrected CNOT gate remain separated by roughly four magnitudes for the wide range of chemical potentials examined.

Outlook. The importance of our results is merely a reflection of the fundamental property of the Majorana modes: Regardless of what model or what specific platform considered, any Majorana mode can be used for braiding and topological quantum computing by diminishing the hybridization error and thus pushing the upper speed limit to large times. Already today, one could easily include negative hybridization for reducing braiding errors in quantum dot based few-site Kitaev chains [28,29], overcoming material and system size constraints for these early systems. Furthermore, we envision that future topological superconductor platforms can be tailored to directly implement negative hybridization or at least make its usage easily accessible. Our work thus suggests that a Majorana-based quantum computer could in principle be improved to reach fault tolerance by exploiting negative hybridization.

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Data availability. The data that support the findings of this article are openly available [44].

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