Quantum Computing with Majorana Kramers Pairs

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We propose a universal gate set acting on a qubit formed by the degenerate ground states of a Coulomb-blockaded time-reversal invariant topological superconductor island with spatially separated Majorana Kramers pairs: the “Majorana Kramers qubit.” All gate operations are implemented by coupling the Majorana Kramers pairs to conventional superconducting leads. Interestingly, in such an all-superconducting device, the energy gap of the leads provides another layer of protection from quasiparticle poisoning independent of the island charging energy. Moreover, the absence of strong magnetic fields—which typically reduce the superconducting gap size of the island—suggests a unique robustness of our qubit to quasiparticle poisoning due to thermal excitations. Consequently, the Majorana Kramers qubit should benefit from prolonged coherence times and may provide an alternative route to a Majorana-based quantum computer.

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In recent years an increasing number of platforms have been proposed for realizing time-reversal invariant topological superconductors (TRI TSCs) [1]. Among the most notable platforms are nanowires and topological insulators in contact to unconventional superconductors (SCs) [2–5] and conventional SCs [6–11], proximity-induced Josephson \(\pi\) junctions in nanowires and topological insulators [12–15] as well as TSCs with an emergent time-reversal symmetry (TRS) [16–19].

A common feature of TRI TSCs is that they host spatially separated Majorana Kramers pairs (MKPs) which form robust, zero energy modes protected by TRS. In spite of much fundamental interest in the properties of MKPs [20–26], a yet unsolved question is if MKPs can be employed for applications in quantum computation. Here, we answer this question in the affirmative.

The purpose of this Letter is to introduce a qubit formed by the degenerate ground states of a Coulomb-blockaded TRI TSC island with spatially separated MKPs: the “Majorana Kramers qubit” (MKQ). We depict the minimal experimental setup for a single MKQ in Fig. 1. It comprises two SC leads which separately couple to two distinct MKPs on a \(U\)-shaped TRI TSC island. The two SC leads are also coupled among themselves by normal and spin-flip tunneling barriers.

Within this setup, we implement single-qubit Clifford gates not by braiding of MKPs [28,29] but rather by making use of a measurement-based approach to Majorana quantum computing [30,31]. In such a measurement-based approach, we perform readout of Majorana bilinears, which correspond to the Pauli operators of our qubit, through Josephson current measurement. More specifically, we selectively deplete either the normal or the spin-flip tunneling barrier and, in this way, realize Josephson couplings that contain different Majorana bilinears, see Fig. 2. To achieve universal quantum computing, we further propose an unprotected \(T\) gate and entangling gate by pulsing of tunnel couplings.

The main conceptual lesson we learn is that Majorana-based quantum computing is possible without the need for magnetic fields. In particular, this point differentiates our proposal from previous works [32–41] that propose quantum computing architectures based on conventional...
Majorana bound states, which in the available experimental candidate platforms require the use of strong magnetic fields for their realization. In particular, the absence of such strong magnetic fields in our proposal constitutes a fundamental advantage because magnetic fields act detrimentally on superconductors and, as a result, constitute a significant challenge toward realizing topological superconductivity. Besides that, there are two features of our setup that yield improved protection from quasiparticle poisoning, (1) within the single-MKP setup of Fig. 1, single-electron tunneling from the SC leads does not only require overcoming the charging energy of the TRI TSC island but also the breaking of a Cooper pair in the leads. Consequently, the SC gap of the leads provides an additional layer of protection against quasiparticle poisoning, independent of the island charging energy. (2) Quasiparticle poisoning due to thermal excitations within the TRI TSC island is suppressed the SC gap of the island itself. In particular, the energy gap of a TRI TSC may be larger than the energy gap of TRS-breaking Majorana islands [32–34] since there is no magnetic field reducing the SC gap size. However, despite the absence of magnetic fields, repulsive interactions, which are present in certain TRI TSC setups [7,13], could suppress the SC gap size.

Setup.—As shown in Fig. 1, our setup comprises a U-shaped TRI TSC islands hosting MKPs $\gamma_{\ell,s}$ with $s = \uparrow, \downarrow$ at spatially well separated boundaries $\ell = L, R$. The two members of a MKP are related by TRS,

$$T \gamma_{\ell,\uparrow} T^{-1} = \gamma_{\ell,\downarrow}, \quad T \gamma_{\ell,\downarrow} T^{-1} = -\gamma_{\ell,\uparrow}.$$  
(1)

We assume that the dimensions of the vertical island segments exceed the MKP localization length $\xi_{\text{MKP}}$ to avoid couplings to fermionic modes that are potentially localized at the island corners [27]. The MKPs are then robust zero-energy states protected by TRS.

Since the TRI TSC island is of mesoscopic size, it acquires (like mesoscopic TSC islands with broken time-reversal symmetry [42]) a charging energy given by

$$U_C = \left(\frac{ne - Q}{e}\right)^2/2C.$$  
(2)

Here, $Q$ is the island gate charges that is continuously tunable with a voltage across a capacitor with capacitance $C$. We tune the gate charge $Q/e$ so that the ground state of the TRI TSC island comprises $n_0$ electron charges. For a sufficiently large charging energy $e^2/2C$, the joint parity of the MKPs on the TRI TSC island is then given by [42,44]

$$\gamma_{L,\uparrow} \gamma_{R,\downarrow} \gamma_{L,\downarrow} \gamma_{R,\uparrow} = (-1)^{n_0}.$$  
(3)

We note that this parity constraint applies to the MKPs on the same TRI TSC, which is different from quantum computing proposals based on conventional Majorana bound states that require two different TSCs that are connected by a conventional SC bridge [39–41]. The parity constraint reduces the fourfold degeneracy of the ground state at zero charging energy, to a twofold degenerate ground state which forms the MKQ. The Pauli operators acting on each of the two MKQs can be written as bilinears in the Majorana operators,

$$\hat{x} = i \gamma_{R,\uparrow} \gamma_{L,\downarrow}, \quad \hat{y} = i \gamma_{R,\downarrow} \gamma_{L,\uparrow}, \quad \hat{z} = i \gamma_{R,\uparrow} \gamma_{L,\downarrow}. $$  
(4)

Under TRS, the Pauli operators transform as $T \hat{x} T^{-1} = (-1)^{n_0} \hat{x}, \quad T \hat{y} T^{-1} = -\hat{y}, \quad$ and $T \hat{z} T^{-1} = (-1)^{n_0} \hat{z}$. We note again that the Pauli operators are defined in terms of MKPs on the same TRI TSC islands, unlike the Pauli operators in quantum computing proposals based on Majorana bound states that are localized on different TSCs connected by a conventional SC bridge [39–41].

In our setup, we choose to address the MKQ by weakly coupling each MKP to a separate $s$-wave SC lead. The Hamiltonian for the two SC leads is of the standard Bardeen-Cooper-Schrieffer form,

$$H_{\text{SC}} = \sum_{\ell=L,R} \sum_{k} \psi_{\ell,k}^\dagger (\xi_{k} \eta_{\ell} + \Delta_{\ell} \eta_{\ell} e^{i\varphi_{k}} \eta_{\ell}) \psi_{\ell,k},$$  
(5)

where $\psi_{\ell,k} = (c_{\ell,k,\uparrow}, c_{\ell,k,\downarrow})^T$ is a Nambu spinor with $c_{\ell,k,s}$ the electron annihilation operator at momentum $k$ and spin $s$ in lead $\ell$. The Pauli matrices $\eta_{\ell,\uparrow,\downarrow}$ act in Nambu space. Furthermore, $\xi_{k}$ is the normal state dispersion and $\Delta_{\ell,\varphi_{k}}$ denote magnitude and phase of the SC order parameter of the $m$-SC lead. The SC phase difference is $\varphi \equiv \varphi_L - \varphi_R$. We assume low temperatures, so no quasiparticle states in

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig2.png}
\caption{(a) A sequence of intermediate states in which a Cooper pair moves between the SC leads by splitting up between the normal tunneling barrier and the TRI TSC. The spin-flip tunneling barrier is fully depleted. The MBS operators (not participating in the sequence are shown in yellow (gray). The resulting effective coupling is \(\tilde{\gamma} = \frac{1}{2} \gamma L \downarrow \gamma L \downarrow\)). (b) Same as (a) but now the Cooper pair splits up between the spin-flip tunneling barrier and the TRI TSC. The normal tunneling barrier is fully depleted. The coupling is \(\gamma L \uparrow \gamma L \downarrow\).}
\end{figure}
the SC leads are occupied with notable probability and can couple to the MKPs.

The most general tunneling Hamiltonian between the MKPs and the fermions on the $\ell$-SC lead reads

$$H_T = \sum_{\ell=L,R} \sum_{k,s} \lambda_\ell c_{\ell,k}^\dagger Y_{\ell,s} e^{-i\phi_k/2} + \text{H.c.}, \quad (6)$$

where we have diagonalized the tunneling Hamiltonian in spin space by an appropriate rotation of the lead fermions [26]. This rotation constraints the pointlike tunneling amplitudes $\lambda_\ell$ to be real numbers, a consequence of the time-reversal symmetry that is not present in quantum computing proposals based on conventional Majorana bound states (MBS) [41]. The operators $e^{\pm i\phi_k/2}$ raise or lower the total island charges by one unit, $[n, e^{\pm i\phi_k/2}] = \pm e^{\pm i\phi_k/2}$, while the MBS operators $Y_{\ell,s}$ flip the respective electron number parities.

As is evident from Fig. 1, there are two types of couplings between the SC leads. The first type is an indirect coupling via the TRI TSC islands which is induced by the tunneling Hamiltonian of Eq. (6). The second type is a direct coupling via two additional tunneling barriers. The first tunneling barrier is used for measuring the $z$-Pauli operator and only allows for normal tunneling,

$$H_N = t_N \sum_k c_{R,k}^\dagger c_{L,k} + c_{L,k}^\dagger c_{R,k} + \text{H.c.}, \quad (7)$$

where $t_N$ is a complex, pointlike tunneling amplitude. The second barrier is used for measuring the $\hat{x}$-Pauli operator and permits spin-flip tunneling,

$$H_S = t_S \sum_k c_{R,k}^\dagger c_{L,k} - c_{L,k}^\dagger c_{R,k} + \text{H.c.}, \quad (8)$$

where $t_S$ is again a complex, pointlike tunneling amplitude. For the $\hat{x}$ ($\hat{z}$) measurement protocols, we require $\text{Im} t_S \neq 0$ ($\text{Im} t_N \neq 0$) when $n_0$ even and $\text{Re} t_S \neq 0$ ($\text{Re} t_N \neq 0$) when $n_0$ odd. In addition, we propose two ways to engineer such tunneling barriers: (1) We consider barriers with a finite intrinsic spin-orbit coupling with spin-orbit length $\lambda_{SO}$ as well as different barrier lengths $d$, $d'$. Tuning $\lambda_{SO}/d'$ ($\lambda_{SO}/d$) to a positive integer (positive half integer) realizes a barrier with pure normal (spin-flip) tunneling [45]. (2) We consider barriers with an engineered spin-orbit coupling due to a local, rotating magnetic field induced by a series of nanomagnets [46,47]. By adjusting the rotating field period through the nanomagnet separation, we can realize barriers with pure normal or spin-flip tunneling.

In summary, the full Hamiltonian reads $H = U_C + H_{SC} + H_T + H_N + H_S$.

**Single-qubit Clifford gates.**—In this section, we will implement single-qubit Clifford gates by “Majorana tracking” [31]. This means for a given circuit of single-qubit Clifford gates we record all Pauli operator redefinitions on a classical computer and use the quantum hardware only to perform suitable measurements of the $\hat{x}$, $\hat{y}$, $\hat{z}$-Pauli operators at the end of the computation.

First, for measuring the $z$-Pauli operator, we consider the situation when a local gate depletes the spin-flip tunneling barrier between the two SC leads, $\text{Im} t_S = 0$ for $n_0$ even and $\text{Re} t_N = 0$ for $n_0$ odd.

In this case, to second order in $t_N$, Cooper pairs tunnel between the SC leads only via the normal tunneling barrier inducing a finite Josephson coupling $J_N \sim 2t_N^2/|\Delta|$. In particular, a Josephson coupling due to Cooper tunneling between each SC lead and the TRI TSC is unfavorable due to the substantial island charging energy [41]. The island charging energy thus plays two key roles. First, it suppresses quasiparticle poisoning due to single-electron tunneling from the environment. Second, it suppresses local mixing $\alpha \hat{y}$ due to Cooper pair tunneling between each SC lead and the TRI TSC island. Such local mixing terms are—as noticed earlier [48]—of importance for TRI TSCs with zero charging energy and, as we will see, can be used to measure the $\hat{y}$-Pauli operator.

Next, we note that to third order in $t_N$, $\lambda_L$, $\lambda_R$ Cooper pair splitting sequences between the TRI TSC island and the normal tunneling barrier induce additional Josephson couplings, $J_z$ for $n_0$ even and $J'_z$ for $n_0$ odd, see Fig. 2(a). In a first process, a Cooper pair on the left SC lead breaks up and one of the electrons tunnels via the normal tunneling barrier to the right SC lead. This leaves the left SC lead in an excited state with one quasiparticle above the SC gap. In a second process, the quasiparticle on the left SC tunnels to the TRI TSC island and increments its charge by one unit. While the left SC returns to its ground state in this way, the TRI TSC island is now in an excited state with one excess charge. It, therefore, requires a third process to remove the extra charge from the TRI TSC by recombining it to a Cooper pair on the right SC lead. Critically, the tunneling events via both the normal tunneling barrier and the TRI TSC island conserve the electron spin. For that reason, the just described third-order sequences contribute terms $\alpha \hat{z} = i\gamma_{R,\downarrow} Y_{L,\downarrow} = (-1)^{n_0} i\gamma_{L,\uparrow} Y_{R,\uparrow}$.

For $\nu \nu_f/\lambda_{SO}^2 \ll \Delta$, $e^2/2C$ with $\nu_f$ the normal-state density of states per spin of the $\ell$-SC lead at the Fermi energy, we compute the amplitudes of all above-mentioned sequences perturbatively. Up to third order in the tunneling amplitudes, we obtain an effective Hamiltonian acting on the ground states of the SC leads and the TRI TSC island.

For $n_0$ even and $n_0$ odd, we find

$$H_{\text{even}} = -(J_N + \hat{z} J'_z) \cos \varphi,$$

$$H_{\text{odd}} = -J_N \cos \varphi + \hat{z} J'_z \sin \varphi. \quad (9)$$

The Josephson couplings are $J_z = \text{Im}(t_N)\nu_f \alpha/\Delta$ and $J'_z = -\text{Re}(t_N)\nu_f \alpha/\Delta$, where we assumed $\lambda_L = \lambda_R \equiv \lambda$, $\nu_L = \nu_R \equiv \nu_f$, and $\alpha$ is a dimensionless prefactor of order one if $U \sim \Delta [49]$. Notably, $J_z$, $J'_z$ and $J$ are of comparable
magnitude if we choose the coupling between the SC leads so that \( \nu_F \lambda^2 \gtrsim 2|t_N|^2/[\Im(t_N) \alpha] \) \( \nu_F \lambda^2 \gtrsim 2|t_N|^2/|\Re(t_N) \alpha| \). Also, we note that both effective Hamiltonians exhibit TRS: for \( H_{z, \text{even}} \) both \( \hat{z} \) and \( \cos \varphi \) are time-reversal even, while for \( H_{z, \text{odd}} \) both \( \hat{z} \) and \( \sin \varphi \) are time-reversal odd.

To measure the \( z \) eigenvalue of the \( \hat{z} \)-Pauli operator, we adopt a two-step protocol. (1) First, we separately measure the Josephson current through the normal tunneling barrier and through the TRI TSC island to determine \( J_N \) and \( J_z \). (2) Second, we measure the Josephson current through the entire device. For \( n_0 \) even, the latter is given by \( I = I_c \sin \varphi \) with the critical current \( I_c = 2e(J_N + zJ_z)/\hbar \) fixing the \( z \) eigenvalue. For \( n_0 \) odd, the current phase relation is of the form \( I = I_c \sin(\varphi + \varphi_0) \). This time it is not the critical current \( I_c = 2esgn(J_N)\sqrt{(J_N)^2 + (J_z)^2}/\hbar \) but the anomalous phase shift \( \varphi_0 = z \arctan(J_z/J_N) \) which fixes the \( z \) eigenvalue. We note that the anomalous phase shift results from the \( \hat{z} \) eigenstates breaking TRS when \( n_0 \) is odd.

To measure the \( x \)-Pauli operator, we consider a depleted normal tunneling barrier, \( \Im t_N = 0 \) for \( n_0 \) even and \( \Re t_N = 0 \) for \( n_0 \) odd. As before, second order cotunneling events now induce a Josephson coupling \( J_s \sim 2|t_s|^2/\Delta \) as a result of Cooper pair tunneling via the spin-flip tunneling barrier whereas fourth order events mediate a Josephson coupling \( J \) via the TRI TSC island. However, a qualitative difference to the preceding considerations arises for the third order Cooper pair splitting sequences, see Fig. 2(b). These sequences now demand two spin flips, one for an electron to move through the spin-flip tunneling barrier and one for an electron to move through the TRI TSC island. Consequently, the third-order sequences now contribute terms \( \propto \hat{x} = i\gamma_{r,s} i\gamma_{L,s} = (-1)^n i\gamma_{r,s} i\gamma_{L,s} \). Up to third order in the tunnel couplings, the effective Hamiltonians for \( n_0 \) even and \( n_0 \) odd read,

\[
H_{x, \text{even}} = -(J_s + \hat{x} J_s) \cos \varphi, \\
H_{x, \text{odd}} = -J_s \cos \varphi + \hat{x} J_s \sin \varphi.
\]

Here, \( J_s = \Im(t_s)\lambda^2 \nu_F \alpha/\Delta \) and \( J_s' = -\Re(t_s)\lambda^2 \nu_F \alpha/\Delta \), where we assumed \( \lambda_L = \lambda_R \equiv \lambda \) and \( \nu_L = \nu_R \equiv \nu_F \). We further note that both effective Hamiltonians exhibit TRS: for \( H_{x, \text{even}} \) both \( \hat{x} \) and \( \cos \varphi \) are time-reversal even, while for \( H_{x, \text{odd}} \) both \( \hat{x} \) and \( \sin \varphi \) are time-reversal odd. To measure the \( \hat{x} \)-Pauli operator, we see that the effective Hamiltonians are of the same form as those in Eq. (9). Hence, our measurement protocol for the \( \hat{x} \)-Pauli operator carries over to \( \hat{x} \)-Pauli operator measurements.

We highlight that potential errors in the \( \hat{x}, \hat{z} \) measurements occur when both \( \Im t_N \neq 0, \Im t_s \neq 0 \) for \( n_0 \) even or \( \Re t_N \neq 0, \Re t_s \neq 0 \) for \( n_0 \) odd. This can happen either if one of the tunneling barriers is not fully depleted, or the barrier lengths \( d, d' \) are not appropriately adjusted to the spin-orbit length \( \lambda_{\text{SO}} \). Fortunately, this constitutes a hardware error addressable prior to experiments. In particular, the error can be made small with a careful design of a conventional Josephson junction.

Finally, we address \( \hat{y} \) measurements. These require tuning the charging energy of the TRI TSC to zero which is attainable—on demand—by coupling the TRI TSC island to a bulk SC through a gate-tunable valve [37]. Critically, even at zero charging energy the value of the joint fermion parity in Eq. (3) remains protected as a result of the lead SC gap. However, unlike in the case of a substantial charging energy, Cooper pairs may now tunnel in a second order process between each SC lead and the TRI TSC island inducing a Josephson coupling \( \propto \hat{y} [20] \). Consequently, the resulting Josephson current provides a means for measuring the \( \hat{y} \) eigenvalue. The details of this measurement scheme are discussed in [20].

**Universal quantum computation.**—For universal quantum computation, the single-qubit Clifford gates need to be supplemented by a \( T = \exp(-i\hat{z}\pi/8) \) gate and an entangling gate \([50]\). If \( n_0 \) odd [even], we obtain the \( T \) gate by pulsing \( J_z \cos(\varphi) [J_x' \cos(\varphi)] \) in \( H_{z, \text{even}} [H_{z, odd}] \) for a duration \( \tau \) so that \( \int_0^\tau J_z(t') \cos(\varphi(t')) dt' = \pi/8 \) \( \{ \int_0^\tau J_x(t') \cos(\varphi(t')) dt' = \pi/8 \} \). Because of imprecisions in the pulsing intervals, these operations are not protected. The need for unprotected gates is generic for Majorana qubits [38–41]. Moreover, for the presented procedure, phase-independent contributions—which were irrelevant for the Josephson current—should now be included in the effective Hamiltonians, see Ref. [49].

For an entangling gate, we consider the setup of Fig. 3 which comprises two SC leads addressing two MKQs \( a, b \). A local gate depletes both normal and spin-flip tunneling barrier, \( t_N = t_s = 0 \). If the width of the SC leads is much smaller than their coherence length \( \xi_{\text{SC}} \), a Cooper pair can split between the two TRI TSC islands and entangle the MKQs. For symmetric couplings and a ground state charge \( n_0 \) for both islands, we have computed the process amplitudes in the weak coupling limit. An anisotropic Heisenberg coupling results,

\[
H_{ab} = J_y \hat{x}_a \hat{y}_b + [J_{xz} + (-1)^{n_0+1} J_{z} \cos \varphi] (\hat{x}_a \hat{z}_b + \hat{z}_a \hat{z}_b).
\]
For the microscopic form of $J_{xz}, J'_{xz}, J_y$, see Ref. [49]. The Heisenberg interaction can be made isotropic by choosing the SC phase difference such that $J \equiv J_y = J_{xz} + (-1)^{n+1} J'_{xz} \cos \phi$. Pulsing the couplings for a time $\tau$ defined by $\int_0^\tau \tilde{J}(\tau')d\tau' = \pi/2$ then implements a SWAP gate via the unitary time evolution operator. The latter, combined with single-qubit gates, allows for universal quantum computing [51].

Conclusions.—We introduced the “Majorana Kramers Qubit” formed by the ground states of a TRI TSC. By coupling a MKQ to SC leads, single-qubit Clifford gates are realized by qubit measurements. A $T$ gate and an entangling gate are realized by pulsing Josephson couplings. The MKQ shows that strong magnetic fields are not needed for Majorana-based quantum computing.

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