

Majorana physics in a Luttinger liquid with attractive interactions

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Majorana zero modes are the hallmark of topological superconductivity. In one-dimensional systems, these zero modes are usually introduced in the context of gapped, mean-field models that do not conserve particle number, such as the Kitaev chain. By nonlocally encoding a conventional fermion across spatially separated Majorana zero modes, these systems become inherently immune to local decoherence. In this work, we show that signatures of Majorana edge physics persist in a number-conserving, gapless Luttinger liquid of spinless fermions with short-range attractive interactions. We identify the two-point correlator as a sharp diagnostic, revealing an edge-to-edge revival whose sign depends on the fermion-number parity. This revival is robust in the thermodynamic limit and persists in the excited states of the system and at different fillings. A simple particle-hole ansatz for the ground state of the system with an odd number of fermions captures the physics of the system for a wide range of interaction strengths, interpolating between the free-fermion limit and the strongly interacting Majorana regime. Finally, we propose a concrete protocol to realize this model with ultracold dipolar molecules or atoms in an optical lattice and to detect the revival via beam-splitter interferometry, opening an experimental route to Majorana physics beyond the conventional gapped-superconductor paradigm.

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Introduction. Majorana fermions are emergent zero-energy modes which are localized at the two edges of a one-dimensional topological superconductor [1,2]. Joined together, these two Majorana modes form a single fermionic degree of freedom which is highly nonlocal and thus robust against local perturbations and decoherence [3]. This robustness makes Majorana fermions a particularly promising platform for topological quantum computation and has motivated an intense search for their realization in both solid-state [4–6] and ultracold atomic systems [7].

The paradigmatic model hosting Majorana fermions is the Kitaev chain, a one-dimensional p -wave superconductor of spinless fermions [1]. The existence of Majorana modes in the Kitaev chain relies on two key ingredients: the model is not number conserving and it possesses a bulk gap which protects the zero modes energetically and justifies the BCS mean-field treatment of the spectrum. In realistic quantum-simulation platforms, however, the microscopic dynamics usually conserves the total number of particles and fermions in one-dimensional systems form a gapless Luttinger liquid, unless the interactions are specifically shaped [8,9].

Whether—and in what form—Majorana physics survives in such conditions is thus a question that is both conceptually and experimentally relevant, recently addressed in number-conserving models [10–15] and in coupled wires [16–25].

In this work, we show that signatures of Majorana edge physics persist even in the conceptually simplest number-conserving, gapless Luttinger liquid of spinless fermions with nearest-neighbor attractive interactions [Fig. 1(a)]:

$$\hat{H} = -t \sum_{i=1}^L (\hat{c}_i^\dagger \hat{c}_{i+1} + \text{H.c.}) + V \sum_{i=1}^L \hat{n}_i \hat{n}_{i+1}, \quad (1)$$

where $\hat{c}_i^\dagger$, $\hat{c}_i$ are fermionic creation and annihilation operators on site i and $\hat{n}_i = \hat{c}_i^\dagger \hat{c}_i$ is the number operator. From now on, we set the hopping amplitude $t = 1$ as the unit of energy. This model is a gapless Luttinger liquid for $V > -2$ and undergoes a phase transition to a phase-separated state for $V < -2$ [8].

In the Luttinger liquid phase with strong attractive interactions, we show that the two-point correlator $\langle \hat{c}_1^\dagger \hat{c}_j \rangle$ evaluated on the ground state of the system decays to zero as j is moved into the bulk of the system, but revives to a finite value when j approaches the opposite edge of the chain [Fig. 1(b)]. The sign of the revival depends on the parity of the total number of fermions in the system. Additionally, we demonstrate that the revival is a robust feature that survives the thermodynamic limit and persists at different fillings and in the low-lying excited states. Furthermore, a particularly simple ansatz proposed in [14] relates the odd and even ground states for a wide range of interaction strengths V .

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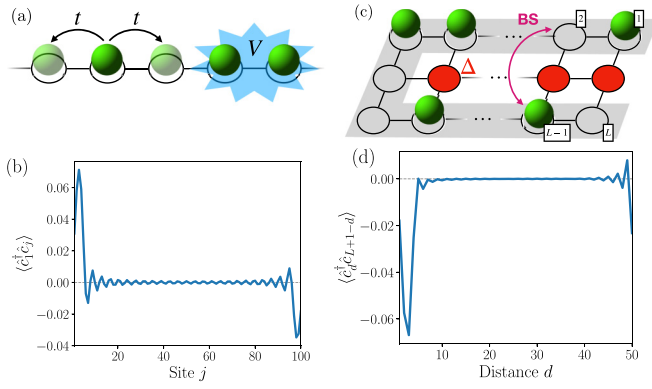


FIG. 1. (a) Schematic representation of the model in Eq. (1). The model describes spin-polarized fermions hopping on a one-dimensional lattice with nearest-neighbor attractive interactions. (b) Two-point correlator $\langle \hat{c}_1^\dagger \hat{c}_j \rangle$ evaluated on the ground state of the model in Eq. (1) at half filling $N = 50, L = 100$ for strong attractive interactions $V = -1.9$. The correlator goes to zero as j is moved into the bulk, but revives to a finite value when j approaches the opposite edge of the chain. (c) Schematic representation of the proposed experimental setup to measure the revival using beam-splitter interferometry. The spin-polarized fermions can hop along a one-dimensional horseshoe (highlighted in gray), since the red sites are blocked by a strong repulsive potential Δ . The labels $1, \dots, L$ of the sites of the chain are written in small black boxes. Performing beam-splitter interferometry between sites d and $L + 1 - d$ allows us to measure the two-point correlator $\langle \hat{c}_d^\dagger \hat{c}_{L+1-d} \rangle$. (d) Two-point correlator $\langle \hat{c}_d^\dagger \hat{c}_{L+1-d} \rangle$ evaluated on the ground state of the model in Eq. (1) at approximately half filling $N = 50, L = 101$ for strong attractive interactions $V = -1.9$. The peak of the revival is located at $d = 3$, which corresponds to a nonzero correlator between two points located at the opposite edges of the chain.

Finally, we propose an experimental realization of the model in Eq. (1) using ultracold dipolar atoms or molecules in an optical lattice [26–31]. The dipoles can hop along a one-dimensional horseshoe-shaped chain [Fig. 1(c)] and the two-point correlator $\langle \hat{c}_d^\dagger \hat{c}_{L+1-d} \rangle$ can be measured using beam-splitter interferometry between sites d and $L + 1 - d$. This correlator exhibits a peak when d corresponds to two points located at the opposite edges of the chain [Fig. 1(d)], which is a smoking-gun signature of the nonlocal nature of the underlying Majorana modes present in the system.

Edge-to-edge revival. In the Kitaev chain, Majorana fermions correspond to two zero-energy modes localized at the left and right edges of the chain [1]. At finite sizes, these two zero modes hybridize and acquire a finite energy splitting [32]. The localization of these two hybridized zero modes can be characterized by the two-point correlator $\langle \hat{c}_i^\dagger \hat{c}_j \rangle$ evaluated on the ground state of the system. For the Kitaev chain, the two-point correlator decays exponentially with the distance $|i - j|$ between the two sites when the first point i is located in the bulk. The presence of modes localized at the edges of the chain is signaled by a revival of the two-point correlator when the first point $i = 1$ is located at one edge of the system. The correlator $\langle \hat{c}_1^\dagger \hat{c}_j \rangle$ goes to zero as j is moved into the bulk of the system, but revives to a finite value when j is located at the opposite edge of the chain. The revival is thus a robust

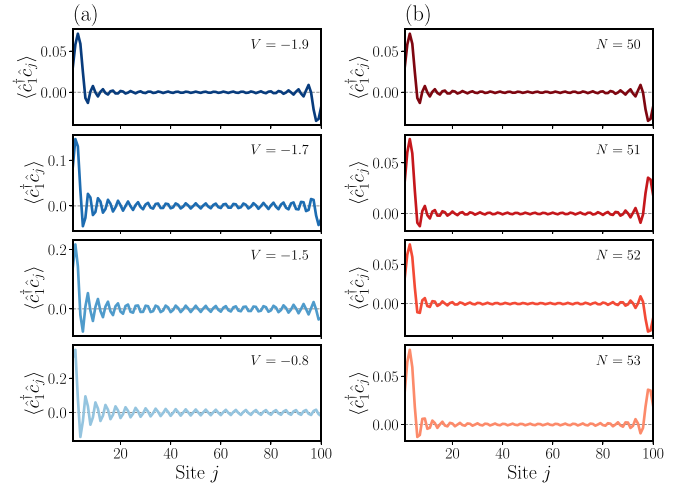


FIG. 2. (a) Two-point correlator $\langle \hat{c}_1^\dagger \hat{c}_j \rangle$ evaluated on the ground state of the model in Eq. (1) at half filling $N = 50, L = 100$ for different values of the interaction strength V . The correlator goes to zero as j is moved into the bulk, but revives to a finite value when j is located at the opposite edge of the chain and V is sufficiently large. (b) Two-point correlator $\langle \hat{c}_1^\dagger \hat{c}_j \rangle$ evaluated on the ground state of the model in Eq. (1) at strong interaction strength $V = -1.9$ and $L = 100$, for different values of the total fermion number N . The revival is robust to adding fermions to the half filled chain, but its sign depends on the parity of the number of fermions N in the system.

fingerprint of the underlying Majorana modes present in the Kitaev model.

In Fig. 2(a), we show that the two-point correlator $\langle \hat{c}_1^\dagger \hat{c}_j \rangle$ evaluated on the ground state of the model in Eq. (1) also exhibits a revival when j approaches the opposite edge of the chain, provided the attractive interactions are strong enough. At the same time, $\langle \hat{c}_i^\dagger \hat{c}_j \rangle \propto |i - j|^{-\alpha}$ decays as a power law in the bulk [8], with an exponent $\alpha = (K + 1/K)/2$ that depends on the Luttinger parameter K of the system.

As a function of the interaction strength V , the revival emerges more distinctly for values of V which approach the critical value $V_c = -2$ from above. For weak attractive interactions, the correlator approaches the free-fermion regime, for which the correlator in the bulk of the system is finite and oscillates with a period of k_F without reviving at the opposite end of the chain [see the lowest panel of Fig. 2(a) for $V = -0.8$]. For stronger attractive interactions, the fermions' tendency to pair is enhanced and the correlator becomes more suppressed in the bulk of the system. This reflects the divergence of the Luttinger liquid parameter, $K \rightarrow \infty$, as V approaches V_c , leading to bulk correlations $\langle \hat{c}_i^\dagger \hat{c}_j \rangle \propto |i - j|^{-K/2}$ becoming arbitrarily short-ranged close to $V_c = -2$.

The model in Eq. (1) conserves the total number of fermions N . In Fig. 2(b), we show the dependence of the revival on the number of fermions N . In particular, we show that the magnitude of the revival is robust to adding fermions to the half filled chain. The sign of the revival, however, depends on the parity of the number of fermions N in the system. This constitutes a direct manifestation of the underlying Majorana modes present in the system, which are known to be sensitive to the parity of the number of fermions.

This can be understood directly for the Kitaev chain at the special point $t = \Delta$ and $\mu = 0$, for which the two Majorana modes are exactly localized at the edges of the chain and they can be expressed in terms of the original fermionic operators $\hat{c}_i$ as $\hat{\gamma}_L = \hat{c}_1^\dagger + \hat{c}_1$ and $\hat{\gamma}_R = i(\hat{c}_L^\dagger - \hat{c}_L)$. The two Majorana operators $\hat{\gamma}_L$ and $\hat{\gamma}_R$ can be used to build a nonlocal fermionic operator $\hat{d}^\dagger = \hat{\gamma}_L - i\hat{\gamma}_R$, which is an exact zero-energy quasiparticle operator. For this reason, the ground state of the Kitaev chain is twofold degenerate and the two ground states $|\text{even}\rangle$ and $|\text{odd}\rangle$ differ by the parity of the number of fermions. In particular, $\hat{d}|\text{even}\rangle = 0$ and the $\hat{d}^\dagger$ quasiparticle operator connects the two ground states with opposite parities

$$\hat{d}^\dagger|\text{even}\rangle = |\text{odd}\rangle. \quad (2)$$

The parity operator $\hat{P}$ in the ground-state subspace is thus

$$\hat{P} = 1 - 2\hat{d}^\dagger\hat{d} = -i\hat{\gamma}_L\hat{\gamma}_R = -(\hat{c}_1^\dagger\hat{c}_L + \hat{c}_1\hat{c}_L + \text{H.c.}). \quad (3)$$

Equation (3) shows that in the Kitaev chain the parity of the ground state $\langle\hat{P}\rangle$ is directly connected to the expectation values of the nonlocal normal and anomalous two-point correlators $\langle\hat{c}_1^\dagger\hat{c}_L\rangle$ and $\langle\hat{c}_1\hat{c}_L\rangle$. In number-conserving models, such as the one in Eq. (1), the anomalous correlator is zero and the sign of the revival $\langle\hat{c}_1^\dagger\hat{c}_L\rangle$ directly dictates the parity of the ground state [see Fig. 2(b)]. More specifically, in the number-conserving case it is possible to build two self-adjoint operators $\hat{\gamma}_L$ and $\hat{\gamma}_R$ (localized at the left and right edges of the chain, respectively) such that the parity $P(N)$ of the ground state $|N\rangle$ with N fermions is given by $P(N) = -ic\langle N|\hat{\gamma}_L\hat{\gamma}_R|N\rangle$, where c is some N -independent constant [14].

It is important to note that the revival is not a finite-size effect [33]. In analogy to the Kitaev chain, the revival is a manifestation of proper, normalized quasiparticle operators, which are unaffected by the $L \rightarrow \infty$ limit. Furthermore, the revival is not a peculiar feature of the ground state of the system at half filling. In fact, we show that the revival is also present at low fillings with a spatial extension which is comparable to the Fermi wavelength of the system and is thus filling dependent [33]. Finally, we find that the revival is also present in the low-lying excited states of the system [33]. In a Luttinger liquid, the low-energy excited states emerge on top of the ground state as long-wavelength density fluctuations, which do not affect the short-wavelength physics of the revival.

Interpolating from Majorana to free fermions. A key signature of the presence of Majorana modes in gapped models that do not conserve particle number, such as the Kitaev chain, is the twofold degeneracy of the ground state. In particular, the two ground states differ by the parity of the number of fermions and they are related by the action of a quasiparticle fermionic operator $\hat{d}^\dagger$ [see Eq. (2)]. In terms of the original fermionic operators $\hat{c}_i$, one can express the quasiparticle operator as

$$\hat{d}^\dagger = \frac{1}{\sqrt{2}} \sum_{j=1}^L (\alpha_j \hat{c}_j^\dagger + \beta_j \hat{c}_j), \quad (4)$$

where the coefficients α_j and β_j are exponentially localized at the edges of the chain and form two orthonormal vectors.

In number-conserving models, such as the one in Eq. (1), a similar construction that relates the even- and odd-parity

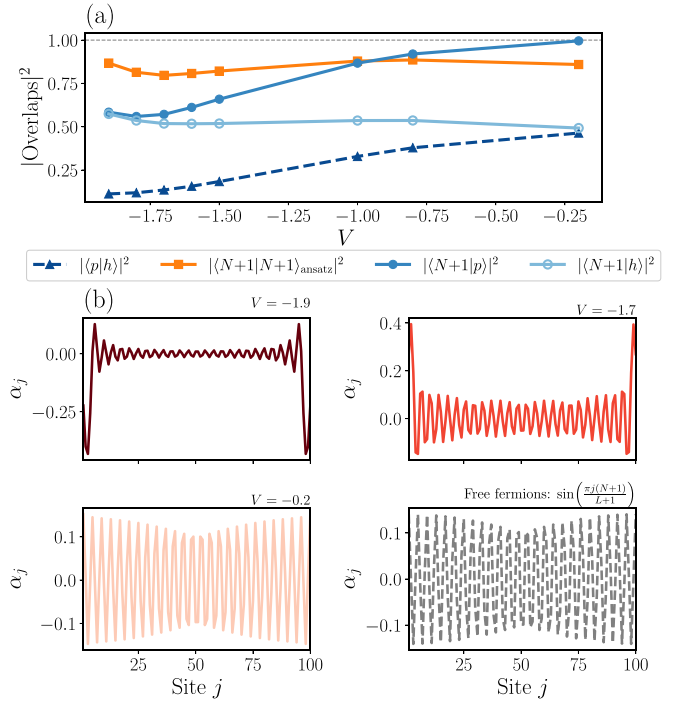


FIG. 3. (a) Squared overlap $|\langle N+1|N+1\rangle_{\text{ansatz}}|^2$ of the ansatz in Eq. (5) with the exact ground state $|N+1\rangle$ of the system with $N+1$ fermions ($N = 50, L = 100$). The squared overlap is very high $|\langle N+1|N+1\rangle_{\text{ansatz}}|^2 \approx 0.85$ for a wide range of interaction strengths V . (b) Coefficients α_i used in the ansatz in Eq. (5) as a function of interaction strength V . The coefficients α_i are localized at the edges of the chain for strong attractive interactions, while they are delocalized in the bulk of the system for weak attractive interactions, smoothly interpolating between the Majorana and free-fermion limits. The analytical expression for the coefficients α_i in the free-fermion limit $V = 0$ is displayed as a gray dashed line.

ground states can be made [14]. In particular, one can make the following ansatz for the ground state $|N+1\rangle$ with an odd number of fermions $N+1$:

$$|N+1\rangle_{\text{ansatz}} \propto \sum_{j=1}^L \alpha_j \hat{c}_j^\dagger |N\rangle + \sum_{j=1}^L \beta_j \hat{c}_j |N+2\rangle. \quad (5)$$

The α_i and β_i vectors are the two orthonormal eigenvectors corresponding to the only nonzero eigenvalues of the matrix $\Delta C_{ij} = \langle N+1|\hat{c}_i^\dagger \hat{c}_j|N+1\rangle - \langle N|\hat{c}_i^\dagger \hat{c}_j|N\rangle$, which can be shown numerically to have approximately rank two [14].

As shown in the left panel of Fig. 3(a), the ansatz in Eq. (5) has a very high squared overlap with the exact ground state of the system $|\langle N+1|N+1\rangle_{\text{ansatz}}|^2 \approx 0.85$ for a wide range of interaction strengths V . To explain this, we decompose the ansatz in Eq. (5) into two parts $|N+1\rangle_{\text{ansatz}} \propto |p\rangle + |h\rangle$, where $|p\rangle \propto \sum_{j=1}^L \alpha_j \hat{c}_j^\dagger |N\rangle$ and $|h\rangle \propto \sum_{j=1}^L \beta_j \hat{c}_j |N+2\rangle$ are the particle and hole components, respectively.

In the limit of strong attractive interactions $V = -1.9$, the particle and hole components are almost orthogonal $|\langle p|h\rangle|^2 \approx 0.1$ and the ground state is well approximated by an equal-weight superposition of the two components $|\langle N+1|p\rangle|^2 \approx |\langle N+1|h\rangle|^2 \approx 0.6$. Furthermore, the α_i and β_i coefficients are localized at the edges of the chain [see

Fig. 3(b)]. This is analogous to the quasiparticle operator $\hat{d}^\dagger$ introduced for the Kitaev chain in Eq. (4), which is an equal-weight superposition of a particle and a hole component.

This picture changes drastically as the interaction strength is decreased, but the ansatz in Eq. (5) still captures the ground state of the system with a high overlap. For $V = -0.2$, the ground state is well approximated by the particle component alone $|\langle N+1|p\rangle|^2 \approx 1$ and the α_i coefficients are delocalized in the bulk of the system. This can be understood explicitly in the free-fermion limit $V = 0$ in an open chain, where the ground state of the system with an odd number of fermions $N+1$ is obtained by adding a fermion with momentum k_F to the ground state of the system with an even number of fermions N . In other words, $|N+1\rangle = \hat{c}_{k_F}^\dagger |N\rangle$, where $\hat{c}_{k_F}^\dagger = \sqrt{2/(L+1)} \sum_{j=1}^L \sin(k_F j) \hat{c}_j^\dagger$ and $k_F = \pi(N+1)/(L+1)$ is the Fermi momentum.

As shown in the right panel of Fig. 3(a), the α_i coefficients interpolate smoothly between the Majorana-fermion regime for strong attractive interactions, in which the coefficients are localized at the edges of the chain, and the free-fermion regime for weak attractive interactions, in which the coefficients are delocalized in the bulk of the system. For this reason, the ansatz in Eq. (5) is a very good approximation of the ground state of the system for a wide range of interaction strengths V .

Experimental realization and outlook. The strong, controllable attractive interactions required to observe the revival in the model in Eq. (1) can be realized with ultracold dipolar quantum gases in an optical lattice. We checked numerically that all the results obtained in this work also apply to the case of dipolar $1/r^3$ interactions, since they are essentially short ranged in one dimension [8]. Two routes are particularly promising: ultracold molecules, whose dipole-dipole interaction is dressed by a microwave field [34–37], giving direct control over both the sign and the strength of the interaction, or atoms with an induced [38] or permanent dipole moment [39–41].

To detect the revival, we propose a single-chain interferometric protocol. The molecules or atoms are loaded in a horseshoe-shaped chain [see Fig. 1(c)] and the ground state of the system is prepared. A beam-splitter operation is then performed between sites d and $L+1-d$, which maps the $\hat{c}_d$ operator to the new operator $\hat{c}'_d$:

$$\hat{c}_d \rightarrow \hat{c}'_d \equiv \frac{1}{\sqrt{2}}(\hat{c}_d + \hat{c}_{L+1-d}). \quad (6)$$

Measuring the occupation number at site d after the beam-splitter operation is equivalent to measuring the number operator

$$\hat{n}'_d = \hat{c}_d^\dagger \hat{c}'_d = \frac{1}{2}(\hat{n}_d + \hat{n}_{L+1-d}) + \frac{1}{2}(\hat{c}_d^\dagger \hat{c}_{L+1-d} + \text{H.c.}), \quad (7)$$

where $\hat{n}_i = \hat{c}_i^\dagger \hat{c}_i$. This protocol, which consists of beam-splitter operations and single-site occupancy measurements, thus allows us to directly reconstruct the two-point correlator $\langle \hat{c}_d^\dagger \hat{c}_{L+1-d} \rangle$ [42] and to detect the revival in the system [see Fig. 1(d)]. Crucially, since the revival survives in the low-lying excited states of the system, a thermal mixture of low-energy

states is sufficient to observe the revival, thus relaxing the cooling requirements for the experiment. Furthermore, an experimentally relevant quantity is given by the ratio of the revival's magnitude to the bulk value of the two-point correlator $|\langle \hat{c}_1^\dagger \hat{c}_{L/2} \rangle|$. Since we are in the Luttinger liquid phase, the bulk value of the two-point correlator near the phase transition $V \rightarrow V_c = -2$ is suppressed as

$$|\langle \hat{c}_1^\dagger \hat{c}_{L/2} \rangle| \propto L^{-K/2} = L^{-1/\sqrt{\epsilon}}, \quad (8)$$

where $\epsilon \propto 2 + V$ at half filling [8]. Remarkably, even in finite-size systems, the two-point bulk correlator $|\langle \hat{c}_1^\dagger \hat{c}_{L/2} \rangle|$ can be heavily suppressed by fine-tuning V close to the critical point V_c . This highlights a major advantage of cold-atom platforms, where the interaction strength V serves as a highly adjustable experimental tuning knob.

An important direction for future work is to understand whether these Majorana modes in spinless Luttinger liquids can be stabilized and protected to realize a topological qubit [43]. Coupling two copies of the model in Eq. (1) to a common third wire, in a spirit analogous to Ref. [16], could similarly pin the edge modes identified in this work. Notably, the specific scaling behavior near the phase transition in Eq. (8) plays a beneficial role by efficiently suppressing single-particle excitations and could provide an advantage in turning this number-conserving, cold-atom-compatible Luttinger liquid into a genuinely topologically protected qubit.

Methods. All the numerical results presented in this work were obtained using DMRG simulations performed with the SYTEN library [44].

Note added. Recently, we became aware of Ref. [14], which originally considered a long-range-interacting model. A discussion of the short-range case appeared in the published version [14].

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Data availability. There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

- [1] A. Kitaev, Unpaired Majorana fermions in quantum wires, *Phys. Usp.* **44**, 131 (2001).
- [2] M. Leijnse and K. Flensberg, Introduction to topological superconductivity and Majorana fermions, *Semicond. Sci. Technol.* **27**, 124003 (2012).
- [3] C. Nayak, S. H. Simon, A. Stern, M. Freedman, and S. Das Sarma, Non-Abelian anyons and topological quantum computation, *Rev. Mod. Phys.* **80**, 1083 (2008).
- [4] R. M. Lutchyn, J. D. Sau, and S. Das Sarma, Majorana fermions and a topological phase transition in semiconductor-superconductor heterostructures, *Phys. Rev. Lett.* **105**, 077001 (2010).
- [5] Y. Oreg, G. Refael, and F. von Oppen, Helical liquids and Majorana bound states in quantum wires, *Phys. Rev. Lett.* **105**, 177002 (2010).
- [6] J. Alicea, New directions in the pursuit of Majorana fermions in solid state systems, *Rep. Prog. Phys.* **75**, 076501 (2012).
- [7] L. Jiang, T. Kitagawa, J. Alicea, A. R. Akhmerov, D. Pekker, G. Refael, J. I. Cirac, E. Demler, M. D. Lukin, and P. Zoller, Majorana fermions in equilibrium and in driven cold-atom quantum wires, *Phys. Rev. Lett.* **106**, 220402 (2011).
- [8] T. Giamarchi, *Quantum Physics in One Dimension*, International Series of Monographs on Physics (Oxford University Press, Oxford, New York, 2003).
- [9] I. Bouchoule, R. Citro, T. Duty, T. Giamarchi, R. G. Hulet, M. Klanjšek, E. Orignac, and B. Weber, Platforms for the realization and characterization of Tomonaga–Luttinger liquids, *Nat. Rev. Phys.* **7**, 565 (2025).
- [10] A. Keselman and E. Berg, Gapless symmetry-protected topological phase of fermions in one dimension, *Phys. Rev. B* **91**, 235309 (2015).
- [11] F. Iemini, L. Mazza, L. Fallani, P. Zoller, R. Fazio, and M. Dalmonte, Majorana quasiparticles protected by $\mathbb{Z}_2$ angular momentum conservation, *Phys. Rev. Lett.* **118**, 200404 (2017).
- [12] J. Ruhman and E. Altman, Topological degeneracy and pairing in a one-dimensional gas of spinless fermions, *Phys. Rev. B* **96**, 085133 (2017).
- [13] S. Fazzini, L. Barbiero, and A. Montorsi, Interaction-induced fractionalization and topological superconductivity in the polar molecules anisotropic $t - J$ model, *Phys. Rev. Lett.* **122**, 106402 (2019).
- [14] J. Thomas-Markarian, K. Agarwal, and I. Martin, Majorana edge modes in isolated wires, *Phys. Rev. Lett.* **137**, 086504 (2026).
- [15] L. Bellinato Giacomelli, T. Bland, L. Lafforgue, F. Ferlaino, M. J. Mark, and L. Barbiero, Topology meets superconductivity in a one-dimensional $t - J$ model of magnetic atoms, *Nat. Commun.* **17**, 5328 (2026).
- [16] L. Fidkowski, R. M. Lutchyn, C. Nayak, and M. P. A. Fisher, Majorana zero modes in one-dimensional quantum wires without long-ranged superconducting order, *Phys. Rev. B* **84**, 195436 (2011).
- [17] M. Cheng and H.-H. Tu, Majorana edge states in interacting two-chain ladders of fermions, *Phys. Rev. B* **84**, 094503 (2011).
- [18] J. D. Sau, B. I. Halperin, K. Flensberg, and S. Das Sarma, Number conserving theory for topologically protected degeneracy in one-dimensional fermions, *Phys. Rev. B* **84**, 144509 (2011).
- [19] C. V. Kraus, M. Dalmonte, M. A. Baranov, A. M. Läuchli, and P. Zoller, Majorana edge states in atomic wires coupled by pair hopping, *Phys. Rev. Lett.* **111**, 173004 (2013).
- [20] F. Iemini, L. Mazza, D. Rossini, R. Fazio, and S. Diehl, Localized Majorana-like modes in a number-conserving setting: An exactly solvable model, *Phys. Rev. Lett.* **115**, 156402 (2015).
- [21] N. Lang and H. P. Büchler, Topological states in a microscopic model of interacting fermions, *Phys. Rev. B* **92**, 041118(R) (2015).
- [22] W.-L. Liu, T.-Z. Yuan, Z. Lin, and W. Yan, Phase diagram of interacting fermionic two-leg ladder with pair hopping, *Chin. Phys. B* **28**, 020303 (2019).
- [23] F. T. Lisandrini and C. Kollath, Majorana edge modes in a spinful-particle conserving model, *Phys. Rev. B* **106**, 245121 (2022).
- [24] N. Tausendpfund, S. Diehl, and M. Rizzi, Majorana zero modes in fermionic wires coupled by Aharonov-Bohm cages, *Phys. Rev. B* **107**, 035124 (2023).
- [25] A. Defossez, L. Vanderstraeten, L. Peralta Gavensky, and N. Goldman, Dynamic realization of Majorana zero modes in a particle-conserving ladder, *Phys. Rev. Res.* **7**, 023183 (2025).
- [26] C. Trefzger, C. Menotti, B. Capogrosso-Sansone, and M. Lewenstein, Ultracold dipolar gases in optical lattices, *J. Phys. B: At. Mol. Opt. Phys.* **44**, 193001 (2011).
- [27] S. Baier, M. J. Mark, D. Petter, K. Aikawa, L. Chomaz, Z. Cai, M. Baranov, P. Zoller, and F. Ferlaino, Extended Bose-Hubbard models with ultracold magnetic atoms, *Science* **352**, 201 (2016).
- [28] J. S. Rosenberg, L. Christakis, E. Guardado-Sanchez, Z. Z. Yan, and W. S. Bakr, Observation of the Hanbury Brown–Twiss effect with ultracold molecules, *Nat. Phys.* **18**, 1062 (2022).
- [29] L. Christakis, J. S. Rosenberg, R. Raj, S. Chi, A. Morningstar, D. A. Huse, Z. Z. Yan, and W. S. Bakr, Probing site-resolved correlations in a spin system of ultracold molecules, *Nature (London)* **614**, 64 (2023).
- [30] L. Su, A. Douglas, M. Szurek, R. Groth, S. F. Ozturk, A. Krahn, A. H. Hébert, G. A. Phelps, S. Ebadi, S. Dickerson, F. Ferlaino, O. Marković, and M. Greiner, Dipolar quantum solids emerging in a Hubbard quantum simulator, *Nature (London)* **622**, 724 (2023).
- [31] A. N. Carroll, H. Hirzler, C. Miller, D. Wellnitz, S. R. Muleady, J. Lin, K. P. ZamarSKI, R. R. W. Wang, J. L. Bohn, A. M. Rey, and J. Ye, Observation of generalized t - J spin dynamics with tunable dipolar interactions, *Science* **388**, 381 (2025).
- [32] M. Cheng, R. M. Lutchyn, V. Galitski, and S. Das Sarma, Splitting of Majorana-fermion modes due to intervortex tunneling in a $p_x + ip_y$ superconductor, *Phys. Rev. Lett.* **103**, 107001 (2009).
- [33] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/9rlr-cqj6> for a detailed numerical analysis of the scaling of the revival, for a discussion of the revival at low and high fillings, and for an analytical treatment using bosonization of the two-point correlator evaluated on the low-lying excited states.
- [34] N. R. Cooper and G. V. Shlyapnikov, Stable topological superfluid phase of ultracold polar fermionic molecules, *Phys. Rev. Lett.* **103**, 155302 (2009).
- [35] T. Karman and J. M. Hutson, Microwave shielding of ultracold polar molecules, *Phys. Rev. Lett.* **121**, 163401 (2018).
- [36] X.-Y. Chen, A. Schindewolf, S. Eppelt, R. Bause, M. Duda, S. Biswas, T. Karman, T. Hilker, I. Bloch, and X.-Y. Luo, Field-linked resonances of polar molecules, *Nature (London)* **614**, 59 (2023).

- [37] N. Bigagli, W. Yuan, S. Zhang, B. Bulatovic, T. Karman, I. Stevenson, and S. Will, Observation of Bose–Einstein condensation of dipolar molecules, *Nature (London)* **631**, 289 (2024).
- [38] E. Guardado-Sanchez, B. M. Spar, P. Schauss, R. Belyansky, J. T. Young, P. Bienias, A. V. Gorshkov, T. Iadecola, and W. S. Bakr, Quench dynamics of a Fermi gas with strong nonlocal interactions, *Phys. Rev. X* **11**, 021036 (2021).
- [39] M. Lu, N. Q. Burdick, and B. L. Lev, Quantum degenerate dipolar Fermi gas, *Phys. Rev. Lett.* **108**, 215301 (2012).
- [40] K. Aikawa, S. Baier, A. Frisch, M. Mark, C. Ravensbergen, and F. Ferlaino, Observation of Fermi surface deformation in a dipolar quantum gas, *Science* **345**, 1484 (2014).
- [41] L. Chomaz, I. Ferrier-Barbut, F. Ferlaino, B. Laburthe-Tolra, B. L. Lev, and T. Pfau, Dipolar physics: A review of experiments with magnetic quantum gases, *Rep. Prog. Phys.* **86**, 026401 (2023).
- [42] A. Impertro, S. Huh, S. Karch, J. F. Wienand, I. Bloch, and M. Aidelsburger, Strongly interacting Meissner phases in large bosonic flux ladders, *Nat. Phys.* **21**, 895 (2025).
- [43] K. Le Hur, Free Majorana fermions with superconducting quantum wires and a magnetic impurity, *C. R. Phys.* **27**, 307 (2026).
- [44] C. Hubig, F. Lachenmaier, N.-O. Linden, T. Reinhard, L. Stenzel, A. Swoboda, M. Grundner, S. Mardazad, and S. Paeckel, The SyTen toolkit, <https://syten.eu/>.