Traveling Majorana Solitons in a Low-Dimensional Spin-Orbit-Coupled Fermi Superfluid

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We investigate traveling solitons of a one- or two-dimensional spin-orbit-coupled Fermi superfluid in both topologically trivial and nontrivial regimes by solving the static and time-dependent Bogoliubov–de Gennes equations. We find a critical velocity $v_h$ for traveling solitons that is much smaller than the value predicted using the Landau criterion due to spin-orbit coupling. Above $v_h$, our time-dependent simulations in harmonic traps indicate that traveling solitons decay by radiating sound waves. In the topological phase, we predict the existence of peculiar Majorana solitons, which host two Majorana fermions and feature a phase jump of $\pi$ across the soliton, irrespective of the velocity of travel. These unusual properties of Majorana solitons may open an alternative way to manipulate Majorana fermions for fault-tolerant topological quantum computations.

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Solitons or localized waves that arise from the interplay between the dispersion and nonlinearity of underlying systems are fascinating phenomena occurring in many different fields of physics [1]. Over the past two decades, a major research emphasis has focused on solitons in atomic Bose-Einstein condensates (BECs) [2]. The family of BEC solitons consists of many interesting members, from bright solitons in attractive BECs [3] and gap solitons in optical lattices [4], to dark solitons in repulsively interacting BECs [5–7], which are created experimentally by imprinting a sharp and characteristic phase jump into the BEC. Remarkably, dark solitons may also be created in strongly interacting Fermi gases [8–16] at the crossover from BECs to Bardeen-Cooper-Schrieffer (BCS) superfluids [17], where phase kinks are encoded in the pairing order parameter. Their recent experimental observation may offer valuable insights into the nature of fermionic superfluidity in the strongly correlated regime [13,18].

Here we consider traveling fermionic solitons in a different setup—one- (1D) or two-dimensional (2D) Fermi superfluids with spin-orbit coupling (see Fig. 1 for a 1D setup)—and predict the existence of an exotic member of the soliton family when the superfluid becomes topologically nontrivial. It is referred to as the Majorana soliton, owing to its ability to host two Majorana fermions that obey non-Abelian statistics at the soliton core [19,20]. Majorana solitons are universal and remarkably robust, in the sense that their properties are not affected by a finite velocity of travel. In particular, the phase jump across a Majorana soliton is exactly pinned to $\pi$ and, therefore, the soliton is not grayed by a finite velocity. This unique stability renders Majorana solitons an ideal platform to manipulate Majorana fermions for topological quantum computations [21].

Our investigation is motivated by the recent realizations of spin-orbit coupling in atomic Fermi gases [22–26] and the promising perspective of creating an atomic topological superfluid [27–29]. Traveling Majorana solitons with the fixed $\pi$-phase step, if experimentally observed to oscillate inside a Fermi cloud, would be a smoking-gun proof of the

FIG. 1. Upper panel: Sketch of the proposed 1D experimental configuration. Lower panel: The magnitude $|\Delta(\xi)|$ and phase $\phi(\xi)$ of the soliton order parameter in the nontopological phase with an interaction parameter $\gamma = 3.41$, Zeeman field $h = 0.52E_F$, and spin-orbit coupling strength $\lambda k_F/E_F = 1.71$. The solid, dashed, and dotted lines correspond to soliton velocities $v_s = 0, 0.15v_F$, and $0.3v_F$, respectively. Here, $k_F$ is the Fermi wave vector and $E_F$ is the Fermi energy.
existence of long-sought topological superfluids. We note that stationary dark solitons with Majorana fermions in a 1D spin-orbit-coupled Fermi gas were recently predicted [30,31]. However, the crucial issue raised in practical manipulations, i.e., the fate of these solitons at a finite velocity of motion, was not addressed.

For concreteness, in the main text we focus on 1D spin-orbit-coupled Fermi superfluids. The results of 2D Fermi superfluids will only be discussed briefly at the end of the Letter, to validate the mean-field theoretical framework used in this work. It is known from previous studies that mean-field theory predicts various qualitative features of 1D interacting quantum gases in the weakly interacting regime. In most cases, the qualitative mean-field predictions, such as the existence of Majorana fermions [32–34] and dark solitons [35], are not invalidated by strong quantum fluctuations in one dimension [36–40]. To ensure the robustness of our 1D mean-field results of Majorana solitons, we carry out additional, extensive simulations for a 2D spin-orbit-coupled Fermi gas [27] and a 2D p-wave superfluid [41], where the use of mean-field theory is widely accepted. The details of these simulations are outlined in the Supplemental Material [42].

Model.—A possible 1D experimental configuration is sketched in the upper panel of Fig. 1. A bundle of parallel, identical 1D spin-1/2 $^{40}$K Fermi gases can be formed by adding a tight 2D optical lattice in the transverse $x$-$z$ plane [33,34], and the spin-orbit coupling with equal Rashba and Dresselhaus weight can be realized by adapting the so-called NIST scheme using two counterpropagating Raman laser beams [22]. The resulting 1D spin-orbit-coupled Fermi gas in a single tube is modeled by the Hamiltonian $H = \int dx [\mathcal{H}_0 + \mathcal{H}_{\text{int}}]$, where [28–31]

$$\mathcal{H}_0 = \left[ \psi^\dagger_\uparrow(x) \psi^\dagger_\downarrow(x) \right] \left[ \left( \mathcal{H}_s + \lambda \hat{k}_s \sigma_y - h \sigma_z \right) \psi_\uparrow(x) \psi_\downarrow(x) \right]$$

(1)

is the spin-orbit-coupled single-particle part and

$$\mathcal{H}_{\text{int}} = g_1 D \psi^\dagger_\uparrow(x) \psi^\dagger_\downarrow(x) \psi_\uparrow(x) \psi_\downarrow(x),$$

(2)

with $g_1 < 0$ is the interaction Hamiltonian describing the attractive contact interaction between the two spin states ($\sigma = \uparrow, \downarrow$). Here, $\psi_\sigma$ is the field operator that creates an atom with mass $m$ in the spin state $\sigma$. The term $\lambda \hat{k}_s \sigma_y - h \sigma_z$ with the momentum operator $\hat{k}_s = -i \partial / \partial x$ and Pauli matrices $\sigma_y$ and $\sigma_z$ is induced by the Rashba process, describing a synthetic spin-orbit coupling with strength $\lambda \equiv \hbar^2 k_R / m$ and an effective Zeeman field $h = \Omega_R / 2$, where $k_R$ and $\Omega_R$ are the momentum and Rabi frequency of the Raman beams [22], respectively. The term $\mathcal{H}_s = -\hbar^2 \partial^2 / (2m) + V_T(x) - \mu$ with the chemical potential $\mu$ describes the motion of atoms in a harmonic potential $V_T(x) = m \omega^2 x^2 / 2$.

We solve the Hamiltonian for stationary and traveling solitons within the mean-field approximation. This amounts to finding solutions with a phase-twisted order parameter in the static and time-dependent Bogoliubov–de Gennes (BdG) equations, $\mathcal{H}_{\text{BdG}} \Phi_\eta(x,t) = E_\eta \Phi_\eta(x,t)$ and $\mathcal{H}_{\text{BdG}} \Phi_\eta(x,t) = i \hbar \left( \partial / \partial t \right) \Phi_\eta(x,t)$, respectively. Here, for convenience, we have used the Nambu spinor representation and have introduced $\Phi_\eta \equiv [u_\eta, v_\eta, v_\eta, u_\eta]$ and $E_\eta$ as the wave function and energy of the Bogoliubov quasiparticles. The BdG Hamiltonian reads

$$\mathcal{H}_{\text{BdG}} \equiv \begin{bmatrix} \mathcal{H}_s - h \lambda \partial / \partial x & 0 & -\Delta \\ \lambda \partial / \partial x & \mathcal{H}_s + h & \Delta \\ -\Delta & 0 & -\lambda \partial / \partial x - \mathcal{H}_s - h \end{bmatrix},$$

(3)

and the BdG equations, either static or time dependent, should be self-consistently solved with the gap equation $\Delta = -g_1 D / \sum |u_\eta|^2 f(E_\eta) + |v_\eta|^2 f(-E_\eta)$ and the number equation $n = -1 / \sum |u_\eta|^2 f(E_\eta) + |v_\eta|^2 f(-E_\eta)$, where $f(E)$ is the Fermi distribution and the sum is performed for the energy level up to a cutoff $E_c$, i.e., $|E_\eta| < E_c$.

To obtain a moving soliton in a trapped gas, we first find a stationary dark soliton at $x_0$ away from the trap center [31]. By evolving such an initial state in time, the soliton is accelerated by the trap potential and caused to oscillate inside the Fermi cloud. The same procedure has previously been used to understand the dynamics of dark solitons in a BEC-BCS Fermi superfluid [9], and could also be employed in experiment.

We also search for traveling soliton solutions on a homogeneous (untrapped) background that satisfy $\Delta(x,t) = \Delta(x - v_s t) = \Delta(x)$, by solving the BdG equations in the comoving frame with the velocity $v_s$ [11]:

$$\mathcal{H}_{\text{BdG}}(\xi) \Phi_\eta(\xi) = [E_\eta - i \hbar v_s \partial / \partial \xi] \Phi_\eta(\xi).$$

(4)

Here, $\mathcal{H}_{\text{BdG}}(\xi)$ is obtained by replacing $\partial_t$ with $\partial_\xi$ and $\Delta(x,t)$ with $\Delta(\xi)$ in Eq. (3). In other words, we seek traveling solitons in a uniform gas that are stationary in the frame of the soliton. This technique provides more insights into the soliton properties and enables us to isolate effects caused by the trapping potential when we analyze time-dependent simulations [12]. For the calculations in a box with length $L$ we impose a modified periodic boundary condition, $\Delta(\xi + L / 2) = \Delta(\xi - L / 2) e^{i \phi}$, to explicitly take into account a phase jump $\delta \phi$ across the soliton [16]. In addition, we implement a generalized secant (Broyden’s) approach to make sure that the self-consistent iteration procedure will converge to a stable solution [11,45].

In our 1D simulations, we use a dimensionless parameter to characterize the interaction strength, $\gamma = -mg_1 D / (\hbar^2 n)$,
which is basically the ratio between the interaction and kinetic energy at the density \( n \). We choose the Fermi vector and energy, \( k_F = \pi n / 2 \) and \( E_F = \hbar^2 k_F^2 / (2m) \), as the units of wave vector and energy, respectively. For simulations in a trapped cloud with \( N \) atoms, it is convenient to use the peak density of a noninteracting Fermi gas in the Thomas-Fermi approximation at the trap center, \( n' = (2/\pi)^{1/2} N m \omega / \hbar \). We denote the corresponding units with \( k_F' \) and \( E_F' \). Throughout this work, we consider only zero temperature. For trapped simulations we shall take the interaction parameter \( \gamma = 3 \), spin-orbit coupling strength \( \lambda k_F / E_F' = 1.5 \), and an energy cutoff \( E_c = 10 E_F' \). Parameters for homogeneous simulations are chosen to correspond to the relevant peak density of the interacting trapped gas.

There are two different regimes for a 1D spin-orbit-coupled Fermi superfluid [27–29,32,33], depending on whether the effective Zeeman field \( h \) is over a threshold \( h_c = \sqrt{\Delta^2 + \mu^2} (\approx E_F \) for the trapped cloud considered). Once \( h > h_c \), the superfluid becomes topologically non-trivial and hosts Majorana solitons. Before presenting our main results on Majorana solitons, it is useful to understand how traveling solitons are affected by spin-orbit coupling in the nontopological phase.

**Nontopological phase.**—The spatial structure of the soliton order parameter in the nontopological phase \((h < h_c)\) is illustrated in the lower panel of Fig. 1. As the velocity increases, the dip in the order parameter profile becomes shallower, and its imaginary part develops structure and becomes larger at the soliton core. Consequently, the phase jump across the soliton decreases from \( \pi \), as shown explicitly in the upper panel of Fig. 2. This turn-to-gray procedure of traveling solitons has been predicted earlier both for BEC-BCS crossover superfluids [11] and BECs [46,47]. However, the presence of spin-orbit coupling leads to some interesting new features.

The most striking feature is that the midgap energy levels of soliton-induced Andreev bound states (ABSs) now exhibit a pronounced velocity dependence, as seen from the lower panel of Fig. 2. Already at zero velocity, the ABS splits into two branches due to the combined effects of spin-orbit coupling and effective Zeeman field [30,31]. With increasing the soliton velocity, the energy of the upper ABS gradually increases and merges into the quasiparticle scattering continuum at \( v_h = 0.22 v_F \), which is much smaller than the pair-breaking velocity \( v_{PB} = 0.45 v_F \). Any coupling between the upper ABS and the bulk continuum states [48] then will destroy the soliton-induced ABSs and in turn make the soliton unstable. Thus, we anticipate that the soliton may decay when its velocity is beyond the threshold \( v_h \), for example, by dissipating its energy in the form of sound waves.

We have checked this conjecture by performing time-dependent simulations in harmonic traps, as reported in Fig. 3. By carefully selecting the position \( x_0 \) of the initially stationary dark soliton, the maximum velocity \( v_m \) reached when the traveling soliton passes the trap center—can be tuned. For \( v_m < v_h \), we find a stable oscillation of the traveling soliton (see the left panel). The oscillation period \( T_s \) seems to satisfy the elegant universal relation (see the inset)

\[
\left( \frac{T_s}{T_x} \right)^2 = \frac{M^*}{M} = 1 + \frac{\hbar n d(\delta \phi)}{2 M^* d v_s},
\]

which was derived by treating soliton as a classical particle [9,47]. Here, \( T_s = 2 \pi / \omega \) is the trapping period, \( M \) and \( M^* \) are, respectively, the physical and inertial mass of the soliton, and their difference is proportional to the derivative of the phase jump [9]. In contrast, at \( v_m > v_h \), the soliton gradually spreads out in the density profile and after a few periods we

![FIG. 2. Upper panel: The phase jump \( \delta \phi / \pi \) as a function of velocity in the nontopological phase. Lower panel: The corresponding mid-gap ABS energy levels. The arrows indicate the velocity \( v_h \), at which the upper ABS (red circles) touches the quasiparticle scattering continuum (shadow area), and the pair-breaking (PB) velocity \( v_{PB} \). Parameters are the same as in Fig. 1.](image1)

![FIG. 3. Time-dependent simulations of traveling solitons in a trapped, nontopological Fermi superfluid with \( h = 0.4 E_F' < h_c \). The color represents the magnitude of density (in units of \( n' \)). By choosing the initial position \( x_0 \), we generate two solitons, whose maximum velocity is \( 0.18 v_F < v_h < 0.22 v_F \) (left panel) and \( 0.42 v_F < v_{PB} < v_h \) (right panel), respectively. The inset examines the universal relation (5) for the soliton oscillation period with or without spin-orbit coupling at different interaction strengths \( (2.5 \leq \gamma \leq 3.2) \). The period \( T_s \) from the time-dependent simulation is compared with \( T_{BdG} \) defined by the right-hand side of Eq. (5) and calculated from the time-independent BdG solutions.](image2)
see only low-amplitude density ripples (right panel). By examining the speed of these ripples, we identify them as sound waves. Our time-dependent simulations with spin-orbit coupling therefore indicate that the critical velocity of traveling solitons could be significantly smaller than Landau critical velocity $v_{PB}$, which was found to be the relevant critical velocity without spin-orbit coupling $[9–11]$. These results are still consistent, since without spin-orbit coupling $v_{h}$ actually is close to the pair-breaking velocity $[12]$.

**Topological phase.**—By increasing effective Zeeman field across $h_{c} = E_{F}$ for a trapped Fermi cloud, the local energy gap (and hence the pair-breaking velocity) at the trap center closes and then reopens. A topological superfluid emerges. The first sign of the existence of a velocity-independent Majorana soliton comes from the time-dependent simulations in harmonic traps, as shown in the upper panel of Fig. 4. During the time evolution, the dip minimum in $\Delta(x,t)$ remains at zero and the phase jump $\delta \phi(t)$ across the soliton is always pinned at $\pi$ (see also the inset in Fig. 5). In the lower panel of Fig. 4, we check more rigorously the velocity dependence using Broyden’s approach. With increasing the soliton velocity in the topological phase, the density and pairing order parameter profiles remain essentially unchanged.

To show the presence of Majorana fermions at the soliton core, we report in Fig. 5 the energy of the ABS as a function of the soliton velocity. Although in the comoving frame the energy $E_{ABS}^{\text{com}}$ increases (linearly) with the velocity, the energy in the laboratory frame, $E_{ABS}^{\text{lab}}$, which is related to the comoving energy by

$$E_{ABS}^{\text{lab}} = E_{ABS}^{\text{com}} + \int d\xi \Phi_{ABS}^*(\xi) \frac{i\hbar \partial \Phi_{ABS}}{\partial \xi},$$

is precisely zero $[49]$. This is expected behavior for a Majorana fermion, which must have zero energy due to the particle-antiparticle symmetry. Together with the observed continuity with the zero velocity case $[30,31]$, we conclude that the moving soliton in the topological phase indeed hosts Majorana fermions. The properties of the Majorana soliton at finite velocity can be made plausible from the universal relation (5), if we assume its validity in the topological phase. We recall that the density notch in Majorana solitons is absent $[30,31]$ and, hence, the physical mass vanishes $[9]$. Equation (5) immediately implies that the derivative of the phase jump is zero, since the oscillation period should be finite. This leads to a constant $\pi$ phase jump, irrespective of the soliton velocity. In turn, the magnitude of order parameter should vanish at the soliton core.

**2D Majorana solitons.**—As shown in the Supplemental Material $[42]$, traveling Majorana solitons exist both in a 2D spin-orbit-coupled Fermi gas in its topological phase and in a 2D $p$-wave Fermi superfluid. In particular, time-dependent simulations in the case of the 2D $p$-wave Fermi superfluid show that the shapes of both density profile and order parameter of the traveling soliton remain the same during the time evolution and thus are independent of the velocity of travel, similar to what happens in the 1D spin-orbit coupled topological Fermi superfluid. All these similarities strongly indicate that the Majorana soliton exists universally, independent of the underlying dimensionality of the topological superfluid.

**Observation of Majorana solitons.**—Experimentally, a 1D or 2D spin-orbit-coupled Fermi gas is easy to set up $[22,43]$, but the realization of topological superfluids is
Majorana solitons can be created by imprinting a $\pi$-phase jump [12,13]. The observation of their oscillations seems to be an experimental challenge, particularly in one dimension, where the density profile of Majorana solitons remains flat. A suitable detection technique is the spatially resolved radio-frequency spectroscopy [50], which may give information about the local order parameter, provided that the size of solitons is comparable with the spatial resolution of spectroscopy. An observation of a vanishing order parameter, provided that the jump [13,15]. The observation of their oscillations seems to be an experimental challenge, particularly in one dimension, where the density profile of Majorana solitons remains flat. A suitable detection technique is the spatially resolved radio-frequency spectroscopy [50], which may give information about the local order parameter, provided that the size of solitons is comparable with the spatial resolution of spectroscopy. 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Majorana solitons in a 2D spin-orbit coupled Fermi gas and in a 2D $p$-wave Fermi superfluid.


[48] Broyden’s approach does not capture such a coupling and can predict traveling soliton solutions at $v_s > v_h$.

[49] Because of the particle-hole symmetry, there is another ABS branch whose energy decreases linearly with the traveling velocity in the comoving frame. In the lab frame, its energy is also precisely zero, giving rise to another Majorana fermion hosted by the Majorana soliton.