The spin-orbit coupling (SOC) plays an important role in many physics in condensed matters \cite{1–4} and ultracold atoms \cite{5–10}, especially for the realization of topological phases and associated protected edge modes \cite{11–14}. In solid materials, the SOC can naturally exist for systems without inversion symmetry \cite{3, 4}. In ultracold atoms, this term can be induced using Raman couplings \cite{7, 15–18}. In recent years, both the one dimensional SOC \cite{19–25} and the two dimensional SOC have been realized in both fermionic \cite{26} and bosonic atoms \cite{27}. These progresses have opened a new avenue in experiments to explore various exotic topological phases \cite{1–4, 28–33}.

In these experiments, the properties of the superfluids depends strongly on the direction of Zeeman field. For out-of-plane Zeeman field, which opens an gap at zero momentum but preserves the inversion symmetry, Bardeen-Cooper-Schrieffer pairing is preferred \cite{34, 35}; however for the in-plane Zeeman field without inversion symmetry, the pairs may carry finite momentum for Fulde-Ferrell-Larkin-Ovchinnikov (FFLO) phase \cite{36–42}. The fate of FFLO phase is still elusive in condensed matter physics and was one of the major concerns in ultracold atoms \cite{39}. The problem is that this phase is very easy to enter the gapless regime due to band tiling effect \cite{43–45}, thus not all FFLO phases can support topological protected Majorana zero modes.

In this work, we propose to realize the fully gapped topological FFLO phase and associated Majorana zero modes \cite{42, 46–48} based on only artificial gauge potential. In this ladder representation the two chain represent two different spins with the same spatial wave function. The vertical arrows and horizontal arrows represent the in-plane and out-of-plane Zeeman field, respectively. The in-plane Zeeman field is induced by a site-dependent rotating Zeeman field. (b) and (c) Typical single particle band structure for $\phi < \pi$ and $\phi > \pi$ (case for inverted band), respectively. The dashed and solid lines represent the cases without and with $h$. Other parameters are $h = 0.5t$, $\Omega = 0.4t$.
should be credible, with which we map out the whole phase diagram. This model is relevant to the recent realized artificial gauge potentials in alkaline and rare-earth atoms with negligible heat effect [49–55].

We consider the following one dimensional model in an optical lattice (see the ladder representation in Fig. 1a),

$$H = H_0 + H_g + V_{\text{int}}. \quad (1)$$

In the first term, we consider the spin-independent hopping between the neighboring sites,

$$H_0 = -\sum_{m,s} \left( ta_{m+1,s}^\dagger a_{m,s} + ta_{m+1,s}^\dagger a_{m,s} + \mu m_{m,s} \right) + h \sum_m \left( a_{m}^\dagger - n_{m\downarrow} \right), \quad (2)$$

where $\phi$ can be viewed as flux per plaquette in Fig. 1a. In an optical lattice, this potential can be realized using different schemes, including laser assisted tunneling [56–58] and driven optical lattice [55, 59, 60]. Very recently, this interaction has also been realized in rare-earth atoms [56–61]. This potential can lead to various applications [63–67]. In these schemes, due to lacking of spontaneous emission in the large detuning limit, the heating effect is negligible. We notice that this potential can be regarded as a site-dependent rotating Zeeman field, which breaks the inversion symmetry and induces spin flipping, thus can play the same role as Rashba SOC, although their mechanisms and realizations are totally different. In following we will show that these two methods can be regarded as “complementary” mechanism to each other.

The phase carried by the Zeeman field can be gauged out by a transformation $a_{m,s}^\dagger \rightarrow a_{m,s}^\dagger e^{-im\phi}$, while the spin-down component is unchanged. Then the above single particle Hamiltonian can be written as

$$\mathcal{H}(k) = H_0 + H_g = \epsilon_k + \mathbf{h}(k) \cdot \sigma, \quad (4)$$

where $\mathbf{h} = (h_z, h_y, h_x)$, with $\sigma = (\sigma_x, \sigma_y, \sigma_z)$ being the Pauli matrices. In above equation, we have

$$\epsilon_k = -t(\cos(k + \phi) + \cos(k)) - \mu, \quad h_x = \Omega, \quad h_y = 0, \quad h_z = h - t(\cos(k + \phi) - \cos(k)). \quad (5)$$

Now the flux phase carried by each plaquette enters the diagonal term, which induces misalignment between the two bands. Two typical band structures ($\phi > \pi$ for inverted band) are presented in Fig. 1b-c. In this model, the momentum dependent terms all appear in the diagonal term and the effective Zeeman field $h_z$, now, is momentum dependent, which, in the presence of in-plane Zeeman fields $h_x$ and $h_y$, will induce spin-momentum locking effect. This is different from the Rashba SOC investigated in literatures [3, 4], in which not $h_z$, but $h_x$ and $h_y$ are odd function of $k$. These two models can not be connected by basis rotation, so this new SOC term can be regarded as a “complementary” mechanism of Rashba SOC. For this reason, the artificial gauge potential can play the same role as the Rashba SOC. Notice that in ultracold atoms, the Rashba coupling strength is determined by the momentum of light, and can not be tuned easily in experiments. By fast modulating the Zeeman field, the effective SOC strength can only be tuned to a weaker value [68–70]. In our model, the “effective” SOC is induced by the cooperation of phase and Zeeman fields, thus can be tuned much easier in experiments.

We next consider the effect of weak attractive interactions on the properties of the superfluids, in which regime the mean field prediction is credible. For the inverted band presented in Fig. 1c, it is possible to choose a proper $\mu$, which occupy the whole lower band and partially the higher band. The tilted band structure means that the mean momentum of the two Fermi points is nonzero. Then the weak attractive interaction between atoms should induce pairings near the two Fermi points, which obviously carry a finite momentum $Q \simeq (k_F^x + k_F^y)$. Notice that the spin-momentum locking effect ensures that in each band the eigenvectors contains only even or odd function of $k$. In this way, the interaction can be rewritten as

$$V_{\text{int}} = -U \sum_{k_{1}, k_{2}, q} Y_{k_{1}, k_{2}, q} a_{k_{1}+k_{2}, q}^\dagger a_{k_{1}+k_{2}, q} a_{k_{1}+k_{2}, q} a_{k_{1}+k_{2}, q} a_{k_{1}+k_{2}, q} \quad (7)$$

where $Y_{k_{1}, k_{2}, q} = a^\dagger (k_{1}+q) a (k_{2}-q) a (k_{2}) a (k_{1})$. We have four different scatterings, which, following the g-ology terminologies [73], are written as: (1) Forward scattering on the same branch, $k_1 \sim k_F$; $k_2 \sim k_F$ (for
$i = 1, 2$, $q \sim 0$, $g_0^{L/R} = U a^a(k_F^i)c^c(k_F^i)c(k_F^i)a(k_F^i)$; (2) Dispersion scattering, $k_1 \sim k_F^i, k_2 \sim k_F^j$, $q \sim 0$, $g_2 = U a^a(k_F^i)c^c(k_F^j)c(k_F^j)a(k_F^i)$; and (3) Backward scattering, $k_1 \sim k_F^i, k_2 \sim k_F^j$, $q \sim k_F^2 + k_F^j$, $g_1 = U a^a(k_F^i)c^c(k_F^j)c(k_F^j)a(k_F^j)$.

In these scatterings, only the dispersion scattering $g_2$ between the right and left movers is important for pairing. The backward scattering $g_1$ is important only near half filling. In Fig. 2a-b, we plot the evolution of $k_F^{1,2}$ and $Q/2$ as a function of $h$ and $\mu$, which exhibit good linearity over a wide range. However their spins change in a non-monotonously with the increasing of momentum $k$ (see Fig. 2c), which is totally different from that in the Rashba SOC model. We plot the evolution of $g$ parameters in Fig. 2d. We find that over a wide range of chemical potential, the scattering is always dominated by the pairing term described by $g_2$ due to nearly opposite spin polarization at the Fermi points. With the increasing of chemical potential, the other three parameters, $g_0^{L}$, $g_0^{R}$ and $g_1$, all decrease to very small amplitude and $g_2 \sim U$. For comparison we have employed the same analysis to the model with Rashba SOC and find that for a sufficient large chemical potential and strong SOC strength, the maximum value of $g_2 \sim U/4$. This result indicates the observation of topological superfluids in our model with much weaker attractive interaction.

In following we underpin the above conclusions via the self-consistent mean-field theory in real space [42]

$$-Un_{m,\uparrow}n_{m,\downarrow} \simeq \Delta_m a_{m,\uparrow}^{\dagger}a_{m,\downarrow} + \text{h.c.} + \frac{|\Delta_m|^2}{U},$$

where $\Delta_m = U\langle c_{m\uparrow}c_{m\downarrow} \rangle$, the pairing strength, are solved self-consistently. We do not consider the Hartree term, which can be included in the chemical potential. This model may support a uniform FFLO phase in some proper parameter regime with (see Fig. 1c),

$$\Delta(x) = \Delta_0 e^{i Q x}.$$
strong coupling regime \((U \sim 4.5 - 5.0t)\). However, it was shown in [74] that with this intermediate or strong coupling, unconventional pairings may become important in some parameter regimes. This result is consistent with the dominated dispersion scattering in Eq. 7.

In this case the corresponding Bogoliubov-de Gennes (BdG) equation can be written as

\[
H_{\text{BdG}} = \begin{pmatrix}
H_0(k) & i\sigma_y \Delta_0 e^{iQx} \\
-i\sigma_y \Delta_0^* e^{-iQx} & -H_0^*(k)
\end{pmatrix},
\]

where \(k = -i\partial_x\). The global phase carried by the order parameter can be gauged out by an unitary transformation, \(U = \text{diag}(e^{iQx/2}, e^{-iQx/2})\), while yields [43, 75],

\[
U^\dagger H_{\text{BdG}} U = \begin{pmatrix}
H_0(k + Q/2) & i\sigma_y \Delta_0 \\
-i\sigma_y \Delta_0^* & -H_0^*(k - Q/2)
\end{pmatrix}.
\]

Since the Zeeman field in the off-diagonal term is independent of \(k\), the momentum shift \(Q/2\) will not enter the off-diagonal term. This is totally different from the model with Rashba SOC, in which \(Q\) will contribute to an effective in-plane Zeeman field and induce a strong tilting effect to the band structures [43]. The topological phase boundary for the above model takes place at \(k = 0\) and \(\pi\) are given by [11],

\[
pf(H(k)\sigma_x) = \Omega^2 - \Delta_0^2 - [h - \mu - 2t \cos(k + \phi)] \times [-h - \mu - 2t \cos(k)], \quad k = 0, \pi.
\]

This index \(\nu = \text{sgn}(\text{pf}(H(k)\sigma_x))\) is used to characterize the topological invariant, in which \(\nu = +1 (-1)\) corresponds to the trivial and topological phase, respectively.

We plot the phase diagram as a function of \(U\) and \(h\) in Fig. 3, focusing on the inverted band case with \(\phi = 7\pi/6\), stimulated by the recent experiments in Refs. [9, 10]. The similar phases can be found for other magnetic flux. We find a large parameter regime for topological FFLO phase (denoted as tFFLO), in which the magnitude of pairing, \(\Delta_m\), is almost uniform as a function of \(h\) and \(\mu\) (see Fig. 3d-e). When \(U = 2.0t\), we find that all the FFLO states are topological nontrivial, which means that even without out-of-plane field \(h\), the system can still be a topological phase. This observation is consistent with Eq. 7, in which not only the gauge potential but also the momentum \(Q\) carried by the Cooper pairs contribute to the topological boundaries. For \(U = 2.8t\), we may have a trivial FFLO phase, which is also fully gapped. This fully gapped phase is attributed to \(Q = k_F^x + k_F^y\) with weak interaction, which ensures that after a parallel move of momentum by \(\pm Q/2\) for the particle and hole Hamiltonian, their Fermi points are exactly coincide in the Fermi surface, at which point an sizable gap is opened by the order parameter. We can even prove exactly that the system is always fully gapped when \(Q\) is determined by the two Fermi points.

It is also possible to observe oscillating FFLO phase (denoted as oFFLO), in which the order parameters oscillate fast in real space due to the competition of two different momenta carried by the Cooper pairs, which can be realized in the small or large chemical potential limit with four Fermi points (see also Fig. 1b). For weak interaction \(U\), the boundaries between tFFLO and oFFLO can be determined by the single particle band structures in Fig. 1b-c. The phase diagram as a function of \(h\) and \(U\) is presented in Fig. 3f, in which the inset shows a finite threshold \(U > U_c \approx 1.8t\) is required to realize pairing. We have numerically verified that this threshold is roughly one half of that in the model with Rashba SOC [42, 75, 76]. For large enough Zeeman field the pairing is completely destroyed in the normal gas (denoted as NG) phase (Fig. 3a-b).

We finally discuss the appearance of Majorana zero energy modes at the two ends in the tFFLO phase regime. In the trivial FFLO phase regime, the system is still fully gapped (see the band structure in Fig. 4c). With the increasing of \(h\), the gap is closed and reopened at \(k = \pi\) (Fig. 4d), and then the edge modes appear in the tFFLO phase. The wave function of the Majorana zero modes are shown in Fig. 4b.

To conclude, in this work we propose to realize topological superfluids and associated Majorana zero modes based on artificial gauge potential, which can be regarded as a “complementary” mechanism to the Rashba SOC. This new method has the advantage of realizing the topological phases with much weaker attractive interaction due to the dominated dispersion scattering near the Fermi surface. We find that all the topological FFLO

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**FIG. 4.** (Color online) Majorana zero modes in an one dimensional chain induced by artificial gauge potential. The phase boundary between trivial FFLO and tFFLO is determined by the appearance of edge modes and \(\nu = -1\) (Eq. 12). Parameters are \(\mu = 0.5t\), \(U = 2.8t\). (b) Wave functions of the edge modes for \(h = 0.4t\), \(\mu = 0.2t\). (c) - (e) Band structures of the BdG equation for \(\mu = 0.5t\) and \(h = 0.1t\) (trivial, \(\nu = +1\)), 0.22t (critical) and 0.5t (tFFLO, \(\nu = -1\)), respectively.
superfluids are fully gapped, and we map out the corresponding phase diagram using self-consistent mean-field theory. Our model has the potential to be the first experimental system to realize the long-sought FFLO phase and Majorana zero modes, in regarding of the above mentioned advantages and the negligible role of heating effect in alkaline and rare-earth atoms for the artificial gauge potentials.

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