

Microscopic Theory of Pair Density Waves in Spin-Orbit Coupled Kondo Lattice

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(Received 11 March 2025; accepted 26 June 2025; published 23 July 2025)

We demonstrate that the discommensuration between the Fermi surfaces of a conduction sea and an underlying spin liquid provides a natural mechanism for the spontaneous formation of pair density waves. Using a recent formulation of the Kondo lattice model that incorporates a Yao Lee spin liquid proposed by the authors, we demonstrate that doping away from half filling induces finite-momentum electron-Majorana pair condensation, resulting in amplitude-modulated pair density waves (PDWs). Our approach provides a precise, analytically tractable pathway for understanding the spontaneous formation of PDWs in higher dimensions and offers a natural mechanism for PDW formation in the absence of Zeeman splitting.

DOI: [10.1103/pdqz-zb8k](https://doi.org/10.1103/pdqz-zb8k)

Introduction—Spatially modulated pair condensates, or pair density waves [1] (PDWs), were originally envisioned by Fulde, Ferrell [2], Larkin, and Ovchinnikov [3] (FFLO) as a high-field superconducting response to the Zeeman splitting of the Fermi surface. The recent observation of PDWs via scanning tunneling microscopy [4] measurements in cuprate [5], iron-based [6,7], and heavy fermion superconductors [8] has sparked a resurgence of interest in this phenomenon. Crucially, in each of these unconventional superconductors the PDW develops without a magnetic field. PDWs naturally form as a response of a uniform pair condensate to a charge or spin density wave (CDW) [1,9]. However, the recent PDW observation in the novel triplet superconductor UTe_2 [8] found that the PDW survives the melting of CDW order [10], suggesting that in this case the PDW is a primary order parameter, developing without the aid of a charge or spin density wave. These observations motivate us to seek a natural mechanism for the spontaneous development of pair density wave order.

The condensation of a PDW requires a pair susceptibility that peaks at a finite momentum. This does not occur naturally in a Bardeen-Cooper-Schrieffer (BCS) superconductor [11], where the logarithmic divergence of the uniform pair susceptibility is cut off by the momentum, or more precisely by $v_F Q$, where v_F is the Fermi velocity and Q the wave vector of the PDW. In a BCS superconductor, a

Zeeman field polarizes the Fermi surfaces, cutting off the zero-momentum instability and favoring a finite-momentum peak in the pair susceptibility, leading to an FFLO-like pair density wave [2,3]. While experiment suggests that pair density waves can spontaneously develop in unconventional superconductors even in the absence of a magnetic field, there is no current consensus on the underlying mechanism.

In this Letter we explore PDW formation in the context of the Kondo lattice model for heavy fermions. The current authors (CPT) have recently proposed a three dimensional, solvable model in which the Kondo screening of a $\mathbb{Z}_2$ spin liquid induces superconductivity in the surrounding conduction sea [12–17]. The underlying spin liquid is a three dimensional generalization of the spin-orbital derivative of the well-known $\mathbb{Z}_2$ Kitaev spin liquid, known as the Yao-Lee spin liquid [18]. At half filling and in the weak coupling limit, Kondo coupling between the spin liquid and the conduction electrons induces odd-frequency triplet superconductivity.

In this Letter, we demonstrate that our model [12] provides a natural mechanism for the emergence of incommensurate pair density waves when the system is doped away from half filling. In particular, its weak coupling description offers a long-sought pathway to understanding the spontaneous formation of pair density waves beyond one dimension [1] in the absence of a field.

At half filling, the model exhibits nesting between electron, hole, and Majorana Fermi surfaces, leading to a logarithmic instability into an electron-Majorana pair condensate under infinitesimal Kondo coupling. Our key result is that doping the system away from half filling, through a shift in the chemical potential, modifies the electron and hole Fermi surfaces while leaving the Majorana Fermi

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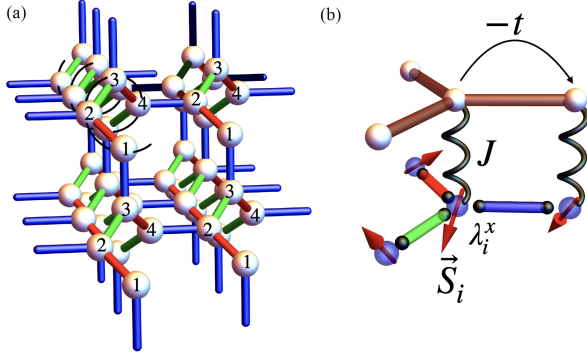


FIG. 1. (a) The hyperoctagon lattice, which features a body-centered cubic structure formed by interpenetrating square and octagonal spirals, with a four-atom unit cell [19]. The frustrated orbital interactions are denoted by color-coded red, blue, and green bonds. (b) In the CPT model each site is trivalent, hosting conduction electrons, shown here linked by golden bonds, localized spins, and orbitals, linked by color-coded orbital interactions. The frustrated orbital interactions induce a Majorana fractionalization of spins, which are then Kondo-coupled to the conduction electrons on the same lattice.

surface unchanged. This imbalance induces a finite-momentum FFLO-like electron-Majorana condensation, giving rise to amplitude-modulated pair density waves.

A key advantage of our model over general parton theories is that it does not rely on an approximate Gutzwiller projection to enforce the single-occupancy constraint by mean-field methods. Instead, in the Yao-Lee spin liquid, this constraint is inherently encoded in the Majorana representation, offering a more precise understanding of how the chemical potential influences the order parameter.

Model—The CPT model [12], defined on the hyperoctagon lattice (Fig. 1), Kondo-couples a band of conduction electrons to the spins of a Yao-Lee spin liquid. Each site possesses 3 degrees of freedom: electrons, localized spins, and localized orbitals, leading to a Hamiltonian with three components, $H_{\text{CPT}} = H_c + H_{\text{YL}} + H_K$, where H_c describes the nearest-neighbor hopping of the conduction electrons, H_{YL} captures the Yao-Lee spin-spin interaction, and H_K couples the conduction sea to the Yao-Lee spin liquid through the Kondo interaction,

$$\begin{aligned} H_c &= -t \sum_{\langle i,j \rangle} (c_{i\sigma}^\dagger c_{j\sigma} + \text{H.c.}) - \mu \sum_i c_{i\sigma}^\dagger c_{i\sigma}, \\ H_{\text{YL}} &= K/2 \sum_{\langle i,j \rangle} \lambda_i^{\alpha_{ij}} \lambda_j^{\alpha_{ij}} (\vec{S}_i \cdot \vec{S}_j), \\ H_K &= J \sum_i (c_i^\dagger \vec{\sigma} c_i) \cdot \vec{S}_i. \end{aligned} \quad (1)$$

The Yao-Lee term involves an anisotropic nearest-neighbor Ising interaction between the $\alpha_{ij} = x, y, z$ components of the orbitals λ_j , which is decorated by a Heisenberg interaction

between the spins $\vec{S}_j$. The orbital frustration in H_{YL} induces large quantum fluctuations, leading to a fractionalization of the spins and orbitals into Majorana fermions as follows, $\vec{\lambda}_j = -i\vec{b}_j \times \vec{b}_j$ and $\vec{S}_j = -(i/2)\vec{\chi}_j \times \vec{\chi}_j$, where b and χ are Majorana fermions. At low energies, the orbital degrees of freedom decouple as static $\mathbb{Z}_2$ gauge fields that freeze into a flux-free configuration at low temperatures, allowing the low energy physics of the Yao-Lee spin liquid to be described by a single-band Hamiltonian,

$$H_{\text{YL}} = \sum_{\mathbf{k} \in \mathbb{B}} \epsilon_\chi(\mathbf{k}) \vec{\chi}_\mathbf{k}^\dagger \cdot \vec{\chi}_\mathbf{k}, \quad (2)$$

where the momentum sum takes place over the cubic Majorana Brillouin zone $\mathbb{B}$ and the dispersion $\epsilon_\chi(\mathbf{k}) = K\tilde{\epsilon}(\mathbf{k})$, where

$$\tilde{\epsilon}(\mathbf{k}) = \frac{1}{2s_x s_y s_z} \left[(c_x^2 + c_y^2 + c_z^2) - \frac{3}{4} \right], \quad (3)$$

where $c_i = \cos(k_i/2)$ and $s_i = \sin(k_i/2)$.

Similarly, the low energy physics of the conduction band [12] can be described by a single band $\epsilon_c(\mathbf{k}) = -t\tilde{\epsilon}(\mathbf{k})$, with $\tilde{\epsilon}(\mathbf{k})$ taking the form (3). By rewriting the conduction Hamiltonian in terms of the Balian-Werthammer spinor $\psi_\mathbf{k} = (c_\mathbf{k}, -i\sigma^2 c_{-\mathbf{k}}^*)^T$, the conduction sea Hamiltonian becomes

$$H_c = \sum_{\mathbf{k} \in \mathbb{B}} \psi_\mathbf{k}^\dagger (\epsilon_c(\mathbf{k}) \mathbb{I} - \mu \tau_3) \psi_\mathbf{k}. \quad (4)$$

The Majorana fractionalization of the spins enables Kondo interaction to be rewritten in terms of Majorana fermions and conduction electrons as follows:

$$H_K = -\frac{J}{2} \sum_j c_j^\dagger (\vec{\chi}_j \cdot \vec{\sigma})^2 c_j. \quad (5)$$

H_K can be decoupled via a Hubbard-Stratonovich transformation in terms of a spinor order parameter field $V_j = -J\langle \vec{\chi}_j \cdot \vec{\sigma} c_j \rangle = (V_{j\uparrow}, V_{j\downarrow})^T$, which is solved self-consistently via a compact mean-field solution for the Kondo interaction,

$$H_{\text{int}}[j] = \frac{1}{2} \left(\psi_j^\dagger (\vec{\sigma} \cdot \vec{\chi}_j) \mathcal{V}_j + \mathcal{V}_j^\dagger (\vec{\sigma} \cdot \vec{\chi}_j) \psi_j \right) + \frac{\mathcal{V}_j^\dagger \mathcal{V}_j}{J}, \quad (6)$$

where $\psi_{\vec{k}} = (c_{\vec{k}}, -i\sigma^2 c_{-\vec{k}}^*)^T$ and $\mathcal{V}_j = (V_j, -i\sigma^2 V_j^*)^T$. In previous work [12], we showed that at half filling, the model exhibits a logarithmic singularity in the electron-Majorana pair susceptibility, leading to odd-frequency triplet superconductivity.

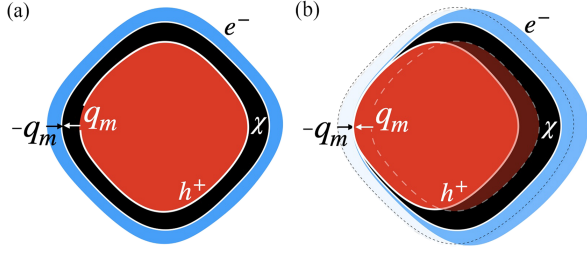


FIG. 2. Illustration of the impact of electron doping on the conduction sea within a hyperoctagon lattice. (a) The conduction sea consists of electron (blue) and hole Fermi surfaces (red), which expand and contract, respectively, as electron doping shifts the system away from half filling. By contrast, the Majorana Fermi surface (black) is unaffected by doping. (b) Doping stabilizes the development of finite-momentum configurations of the electron-Majorana spinor order by shifting the electron and hole Fermi surfaces by $\pm q_m$ to nest with the Majorana Fermi surface.

Pair density waves—We now show that doping provides a natural mechanism for the spontaneous formation of a PDW with a modulated gap function, in which the chemical potential plays a role analogous to a Zeeman field in conventional superconductors. The introduction of chemical potential splits the electron and hole Fermi surfaces while leaving the Majorana Fermi surface unaffected [Fig. 2(a)]. This then causes the spinor order parameter $\mathcal{V}_{\vec{q}}$ to acquire a finite-momentum $\vec{q}$ [Fig. 2(b)], resulting in the formation of a PDW. The Fourier transform of the Kondo interaction with finite-momentum order parameter is given by

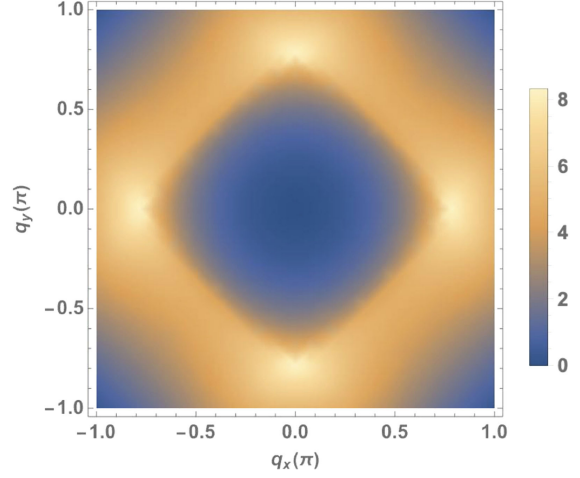


FIG. 3. Depicts the finite-momentum susceptibility $\vec{q} = \vec{Q} - \vec{P}$ for the formation of electron-Majorana spinor order $\mathcal{V}_{\vec{q}}$ along the $q_z = 0$ plane. The susceptibility reaches its maximum at $\vec{q}_i = q_m(\pm 1, 0, 0), q_m(0, \pm 1, 0)$, where $q_m = 0.7738\pi\mu/t$, illustrating the fourfold degeneracy. Because of cubic symmetry, the susceptibility is also maximized at $\vec{q}_i \in \{q_m(\pm 1, 0, 0), q_m(0, \pm 1, 0), q_m(0, 0, \pm 1)\}$.

$$H_{int} = \sum_{\vec{k} \in \mathbb{D}} \left(\psi_{\vec{k}+\vec{q}}^\dagger (\vec{\sigma} \cdot \vec{\chi}_{\vec{k}}) \mathcal{V}_{\vec{q}} + h.c. \right) + \mathcal{N} \frac{\mathcal{V}_{\vec{q}}^\dagger \mathcal{V}_{\vec{q}}}{J}. \quad (7)$$

Finite-momentum spinor order arises when the electron-Majorana pair susceptibility (Fig. 3)

$$\frac{\partial^2 F_e}{\partial V^2} = \chi_P(\vec{q}) = -2T \sum_{i\omega_n, \vec{k} \in \mathbb{D}} \text{Tr} \left[G_\psi(\vec{k} + \vec{q}, i\omega_n) \sigma^a \mathcal{Z} G_{\chi^a}(\vec{k}, i\omega_n) \mathcal{Z}^\dagger \sigma^a \right], \quad (8)$$

where F_e is the electronic part of the free energy, develops a maximum at a finite-momentum $\vec{q}$. Here, $G_\psi(\vec{k} + \vec{q}, i\omega_n)$ and $G_{\chi^a}(\vec{k}, i\omega_n)$ are the conduction and Majorana Green's functions, respectively, with $\mathcal{Z} = (\mathcal{V}/\sqrt{2}V)$ representing the normalized orientation of the spinor order ($\mathcal{Z}^\dagger \mathcal{Z} = 1$). Further simplifications, outlined in Supplemental Material [20], allow us to rewrite the susceptibility for the hyperoctagon lattice in the following closed form,

$$\chi_P(\vec{q}) = \frac{3\rho(0)}{(K+t)} \ln \left(\frac{(K+t)D}{\mu} \right) - \frac{3\rho(0)}{4(K+t)} \int \frac{d\Omega}{4\pi} \left[\ln \left(\left| 1 - \frac{\vec{v}_F(\theta, \phi) \cdot \vec{q}}{\mu} \right| \right) + \ln \left(\left| \frac{K}{t} + \frac{\vec{v}_F(\theta, \phi) \cdot \vec{q}}{\mu} \right| \right) \right]. \quad (9)$$

Here, $\rho(0)$ is the density of states at the Fermi surface, and $\vec{v}_F = \vec{\nabla}_{\vec{k}} \epsilon_c(\vec{k})$ is the Fermi velocity of the conduction sea, which, for the hyperoctagon lattice, depends on the orientation (θ, ϕ) . The logarithmic instability of the electron-Majorana susceptibility in Eq. (9) is cut off by the chemical potential, which splits the nesting between the

conduction sea and the Majorana Fermi surfaces. This results in the susceptibility maxima developing at a finite momenta $\vec{q}$ (Fig. 3).

To evaluate the susceptibility in angular coordinates, we rewrite the Fermi velocity of the conduction sea in the momentum basis,

$$\vec{v}_F = \frac{1}{2s_x s_y s_z} \left(\frac{c_x^3}{s_x}, \frac{c_y^3}{s_y}, \frac{c_z^3}{s_z} \right), \quad (10)$$

where $c_i = \cos(k_i/2)$ and $s_i = \sin(k_i/2)$. We then parametrize the c_i of the Fermi surface in spherical coordinates,

$$(c_x, c_y, c_z) = [r \cos(\phi) \sin(\theta), r \sin(\phi) \sin(\theta), r \cos(\theta)], \quad (11)$$

in terms of angular variables θ and ϕ and radius $r = \sqrt{3}/2$. Carrying out the angular integral in Eq. (9), we numerically find that the susceptibility is maximized (Fig. 3) for

$$\vec{q}_i \in q_m(\pm 1, 0, 0), \quad q_m(0, \pm 1, 0), \quad q_m(0, 0, \pm 1) \quad (12)$$

along the principal axes where the magnitude of the wave number at the susceptibility maxima $q_m = 0.77\pi\mu$ (Fig. 3) is proportional to the chemical potential.

At zero temperature, the pairing susceptibility $\chi_P(\vec{q}_m)$ peaks at finite momentum of magnitude $|\vec{q}_m| \propto \mu$, favoring finite-momentum order. At the superconducting transition temperature T_c , staggered order develops provided $\chi(T, \vec{q})$ is maximal at a nonzero momentum $\vec{q}_m$, implying convexity of the pair susceptibility, $b(T, \mu) = d^2\chi(T, q)/dq^2|_{q=0} > 0$ at $\vec{q} = 0$. Detailed calculation shows that $b(T, \mu) > 0$ is positive provided the chemical potential exceeds $\mu_c = 3.82T$ (see Supplemental Material [20]). Beyond this point, the system undergoes a second-order transition from the normal state into an incommensurate PDW upon cooling. In the absence of fluctuation, mean-field theory predicts that the uniform state at $\mu < \mu_c$ and the PDW at $\mu > \mu_c$ are separated by a first-order transition (Fig. 4). The role of long-wavelength amplitude

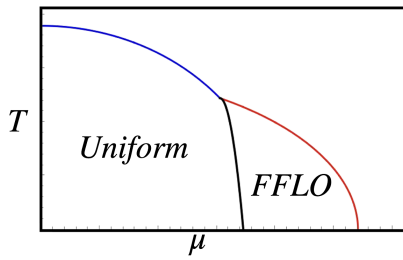


FIG. 4. Depiction of the phase diagram of the CPT model as a function of temperature T , and chemical potential μ for a finite Kondo coupling J . Second-order transitions separate the normal state from the uniform superconducting (blue) and incommensurate FFLO-like PDW phases (red). The first-order boundary (black) between the two ordered phases appears where the critical temperatures coincide. This black line terminates at a critical chemical potential μ_c , marking a first-order transition between the two phases at zero temperature.

fluctuations in shaping this phase boundary is left for a future work.

In the PDW phase, the cubic symmetry of our model leads to six degenerate wave vectors along which the spinor order prefers to modulate. Such degeneracy is often lifted, as in the case of FFLO [3] states, where the order parameter favors amplitude modulation due to attractive interactions between $\mathcal{V}_{\vec{q}}$ and $\mathcal{V}_{-\vec{q}}$. This results in ordering of the form $\mathcal{V}_{\vec{q}}e^{i\vec{q}\cdot\vec{x}} + \mathcal{V}_{-\vec{q}}e^{-i\vec{q}\cdot\vec{x}}$, or other superpositions, depending on the level of doping.

Landau theory—In the incommensurate phase there will in general be up to six modulation vectors q_i ($i = 1, 6$). The Landau theory describing the energetics of this state depends on the six spinors $\mathcal{V}_{q_i}$ and has the general form

$$\begin{aligned} \mathcal{F} = & \sum_i \left[\tau(\mathcal{V}_{q_i}^\dagger \mathcal{V}_{q_i}) + \frac{\beta_1}{2} (\mathcal{V}_{q_i}^\dagger \mathcal{V}_{q_i})^2 \right] \\ & + \sum_{i \neq j} \frac{\beta_2}{2} (\mathcal{V}_{q_i}^\dagger \mathcal{V}_{q_i})(\mathcal{V}_{q_j}^\dagger \mathcal{V}_{q_j}) + \frac{\beta_3}{2} (\mathcal{V}_{q_i}^\dagger \mathcal{V}_{q_j})(\mathcal{V}_{q_j}^\dagger \mathcal{V}_{q_i}) \\ & + \frac{\beta_4}{2} (\mathcal{V}_{q_i}^\dagger \mathcal{V}_{q_j})(\mathcal{V}_{-q_i}^\dagger \mathcal{V}_{-q_j}). \end{aligned} \quad (13)$$

Minimizing (13) with respect to $\mathcal{V}_{q_i}$ determines the stable configurations of the PDW.

The quadratic degeneracy among the order parameters $\mathcal{V}_{q_i}$ is broken by the quartic coefficients β_i , allowing for solutions in which multiple $\mathcal{V}_{q_i}$ acquire an expectation value. Let us consider the simplest PDW with unidirectional modulation in which the $\mathcal{V}_{q_i} = V\mathcal{Z}_{q_i}$ ($i = 1, 2$). In general, we have to allow for the possibility that the two spinors have a different orientation, so that $\mathcal{Z}_{q_i}^\dagger \mathcal{Z}_{-q_i} = \cos^2\phi/2$ is not equal to unity. The modulated phase becomes stable relative to the unmodulated case provided $\beta_2 + \beta_3 \cos^2\phi/2 < \beta_1$; moreover, if $\beta_3 > 0$, then $\phi \neq 0$ and the two spinors in the PDW are no longer aligned. We will shortly see that even when $\phi = \pi$ and the spinors are orthogonal, the electron self-energy is still modulated. More complex intertwined order is possible, giving rise in general to a multi-wave-vector PDW,

$$\mathcal{V}(\vec{x}) = V_0 \sum_i \mathcal{Z}_i e^{i\vec{q}_i \cdot \vec{x}}, \quad \mathcal{Z}_i^\dagger \mathcal{Z}_i = 1. \quad (14)$$

Unlike standard models [2,3], this triplet PDW is described in terms of the six independent spinors $\mathcal{Z}_i$ enabling a rich variety of modulated orders.

Modulated pairing and intertwined order—The modulation in the electronic response that arises from the modulated spinor ordering is encoded in the electronic self-energy, obtained by integrating out the Majorana fermions,

$$\Sigma(x, x', \omega) = \sigma_a \mathcal{V}(x) G_\chi(x - x', \omega) \mathcal{V}^\dagger(x') \sigma_a. \quad (15)$$

The local component of this self-energy, obtained by setting $x = x'$, can be utilized to study the modulated order resulting from the PDW (14).

To illustrate this point, we consider the simplest PDW with unidirectional modulation, which involves two condensed spinors so that $\mathcal{Z}_i, \mathcal{Z}_{-i} \neq 0$. Depending on the coefficient β_3 of the Landau theory, $\mathcal{Z}_i$ and $\mathcal{Z}_{-i}$ may not be equal. The various intertwined components of the modulated order can be understood by decomposing the self-energy,

$$\Sigma(x, \omega) = \Sigma_{\alpha\beta}(x, \omega) \sigma^\alpha \tau^\beta, \quad (\alpha, \beta \in 0, 3), \quad (16)$$

in terms of the spins and isospin operators σ^α and τ^β , respectively.

The components of this four dimensional matrix describe an intertwined order between the primary, modulated triplet pairing and secondary charge, singlet pair, and spin density waves. For example, the spacelike components $\Sigma_{ab}(x, \omega)$ ($a, b = 1, 2, 3$) define the primary, modulated triplet order with a d vector $[\vec{d}^b]_a(x, \omega) = \Sigma_{ab}(x, \omega) \propto \cos(2\vec{q} \cdot \vec{x})$.

By contrast, the mixed components $\Sigma_{0a}(\vec{x}, \omega)$ and $\Sigma_{a0}(\vec{x}, \omega)$ describe secondary charge, singlet pair, and spin density wave components. Terms with only a single σ^a or τ^b pick up a minus sign under transposition, which causes these components to be proportional to $\sin(2\vec{q} \cdot \vec{x})$, so the secondary order is $\pi/2$ out of phase with the primary order (see Supplemental Material [20] for details). Thus $\Sigma_{03}(\vec{x}, \omega) \sim \sin(2\vec{q} \cdot \vec{x})$ describes a staggered scattering potential, while $\Delta(\vec{x}, \omega) = \Sigma_{01}(\vec{x}, \omega) + i\Sigma_{02}(\vec{x}, \omega) \sim \sin(2\vec{q} \cdot \vec{x})$ describes a modulated s -wave pairing, resulting from the absence of inversion symmetry. Similarly, the spin density wave component $\Sigma_{a0}(x, \omega) = \vec{M}_a(\omega) \sin(2\vec{q} \cdot \vec{x})$ results from the time-reversal symmetry breaking of the spinor order. We note that when the two spinors $\mathcal{Z}_{\pm q}$ are aligned, only the triplet and resonant scattering components of the self-energy, i.e., $\Sigma_{\mu\mu}$, remain finite. The appearance of the same modulation vector in the out-of-phase pair density wave and charge density wave is a unique feature of this PDW mechanism.

Discussion—In this Letter, we have proposed a mechanism for the spontaneous formation of pair density waves, where the underlying superconductivity arises from pairing with a $\mathbb{Z}_2$ spin liquid. Screening of the spin liquid by the conduction sea then gives rise to triplet pair density waves. While our model represents a specific case, the underlying concept extends to other forms of unconventional superconductivity believed to result from spin fractionalization, such as the resonating valence bond theory of high- T_c superconductivity.

Pair density wave order has indeed been observed in STM experiments in the heavy fermion superconductor UTe_2 [8], where it has been identified as the parent order underlying secondary charge density wave order, through melting transitions [10]. Furthermore, time-reversal

symmetry breaking observed in chiral step-edge experiments [21] and the onset of reentrant superconductivity at 40 T [22], alongside the absence of a two-stage transition and the fully gapped spectrum revealed by transport measurements [23,24], have fueled debate [21,23–43] over the nature of superconducting order of UTe_2 . This controversy arises because there are no known two-dimensional representations of superconducting order in orthorhombic (Immm) UTe_2 , available to reconcile these experimental findings.

While our Letter does not provide a microscopic mechanism for the superconductivity in UTe_2 , it points to a novel way out of this dilemma. Superconductivity breaks the $U(1)$ isospin symmetry associated with charge conservation. In principle, group theory allows for both singly connected representations of the order parameter, exemplified by the charge $2e$ BCS pairing, and “double-group” representations of the order parameter, with half-integer isospin. Majorana-mediated superconductivity is a first concrete example of this second class of broken symmetry. We may think of the electron-Majorana spinor as analogous to the quark in strong interaction physics, and using this analogy, new families of superconducting phases may then be possible, generated by product group representations akin to the eightfold way of mesons. In such a zoo of enriched superconducting phases, some may be topological [44]. Importantly, double groups introduce Kramers degeneracy, allowing for broken time reversal in low-symmetry crystalline environments, providing an alternative path for the emergence of spontaneous time-reversal symmetry breaking from a heavy Fermi liquid [45].

Acknowledgments—A. P. and P. C. would like to thank Tamaghna Hazra for fruitful discussions. This work was supported by Office of Basic Energy Sciences, Material Sciences and Engineering Division, U.S. Department of Energy (DOE) under Contracts No. DE-SC0012704 (AMT) and No. DE-FG02-99ER45790 (A. P. and P. C.).

Data availability—No data were created or analyzed in this Letter.

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