Symmetry breaking via Kondo hybridization: Chiral nematic metal in Pr$_2$Ir$_2$O$_7$

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In frustrated magnets when magnetic ordering is suppressed down to low temperature, the formation of a quantum spin liquid becomes a possibility. How such a spin liquid manifests in the presence of conduction electrons is a question with potentially rich physical consequences, particularly when both the localized spins and conduction electrons reside on frustrated lattices. We propose a novel mechanism for symmetry breaking in systems where conduction electrons hybridize with a quantum spin liquid through Kondo couplings. We apply this to the pyrochlore iridate Pr$_2$Ir$_2$O$_7$, which exhibits an anomalous Hall effect without clear indications of magnetic order. We show that Kondo hybridization between the localized Pr pseudo-spins and Ir conduction electrons breaks some of the spatial symmetries, in addition to time-reversal regardless of the form of the coupling. These broken symmetries result in an anomalous Hall conductivity and induce small magnetic, quadrupolar and charge orderings. Further experimental signatures are proposed.

**Introduction:** The study of interactions between itinerant electrons and localized degrees of freedom has lead to an understanding of a wealth of novel physical phenomena. These range from isolated moments, as in the Kondo effect\([1, 2]\) through into the realm of dense lattices of moments as in heavy fermion materials\([3–5]\) and the anomalous Hall effect (AHE)\([6]\). While still largely unexplored, the interplay between itinerant degrees of freedom and frustrated local moments promises to unveil new and unique phases of matter\([7]\).

One particularly interesting scenario arises when the local moments are highly frustrated, realizing a quantum spin liquid. How such a spin liquid competes with Kondo hybridization when conduction electrons are present has yet to be fully addressed\([8–11]\).

In this letter, we study systems where conduction electrons interact with a quantum spin liquid, introducing a novel mechanism for breaking spatial symmetries. When the conduction electrons hybridize with spinons the emergent gauge structure of the spin liquid is exposed. We propose that a spin liquid with non-trivial gauge structure, i.e. fluxes penetrating the lattice, is incompatible with trivial gauge structure in the conduction states as well as some of the spatial symmetries. We apply this to a model of conduction electrons and local moments on the pyrochlore lattice, where the effective fluxes are provided by choosing local quantization axes for the conduction electrons. While this is simply a basis choice when the electrons are isolated, when coupled with a fully symmetric U(1) spin liquid on the local moments any uniform hybridization forces the emergent magnetic flux through the plaquettes between the local moments and the conduction electrons (as shown in Fig. 1b) breaking some of the spatial symmetries.

A puzzling example of a material with frustrated local moments and conduction electrons arises in the pyrochlore iridate Pr$_2$Ir$_2$O$_7$, where the Praeseodymium (Pr) and Iridium (Ir) atoms form a pair of interpenetrating pyrochlore lattices with space group $Fd\bar{3}m$, as shown in Fig. 1a. The lack of indications of magnetic ordering\([12]\) well below the Curie-Weiss temperature\([13–14]\) suggests that the Pr sub-lattice is frustrated, either intrinsically or due the presence of the Ir conduction electrons\([15]\). This is corroborated by features in the field dependent magnetization at low temperatures suggesting an anti-ferromagnetic interaction and possibly spin-ice physics, in contrast to the sign of the Curie-Weiss temperature\([16]\). In addition to these magnetic features, the compound is metallic\([13]\), originating in the Ir sublattice, and shows a finite AHE at intermediate temperatures between $\sim 0.3K$ and $\sim 1.5K\([16]\$. The presence of an AHE along [111] indicates a breaking of time-reversal symmetry as well the rotational symmetry of the lattice.

Generically, since magnetization and the anomalous Hall vector $\mathbf{\sigma}_A = \sigma_{xy} \hat{x} + \sigma_{xz} \hat{y} + \sigma_{yz} \hat{z}$ transform in an identical fashion one expects the two orderings to appear together, as is found in ferromagnets\([6]\). The mystery in Pr$_2$Ir$_2$O$_7$ is that the intermediate phase shows no evidence for net magnetiza-
tion to a resolution of $\sim 10^{-3} \mu_B/Pr$ [16]. A number of other unconventional features, such as the lack of a clear phase transition into this intermediate phase as well as unusual behaviour of the Hall conductivity in large fields [13, 14, 16], further enrich the problem. These unexplained properties have attracted considerable theoretical attention, with proposals exploring the full range of scenarios from the interplay between spin-ice physics and the conduction electrons [16, 17] to detailed considerations of the Ir physics [18, 19] and Pr-Ir couplings [20, 21]. While a consensus has yet to emerge, it is clear that both the Pr and Ir degrees of freedom must be taken into account to explain the fascinating phenomena seen in experiments. We employ ideas to understand Pr$_2$Ir$_2$O$_7$, finding that a uniform U(1) spin liquid is favoured, leading to breaking of the appropriate symmetries to allow an AHE when hybridization is included. The orbital nature of this symmetry breaking provides a simple explanation for both the AHE as well as the smallness of the induced magnetic and quadrupolar moments.

**Conduction electrons:** We first construct a minimal model for Pr$_2$Ir$_2$O$_7$, beginning with the Ir atoms. Assuming an ionic configuration of Ir$^{4+}$ one has five $d$ electrons per Ir. These Ir$^{4+}$ ions form a pyrochlore lattice, face centered cubic with a tetrahedral basis, each surrounded by oxygens. Due to the strong octahedral crystal fields and spin-orbit coupling, one can consider only a single half-filled $j_{\text{eff}} = 1/2$ band [22]. In the global cubic axes a symmetry operation $S$ rotates the spin and orbital degrees of freedom according to some representation $R_S$. How these symmetry operations act with the local axes can be seen most clearly if we adopt quantization axes for the $j_{\text{eff}} = 1/2$ states that are compatible with the exact $D_{4d}$ site symmetry of the Ir$^{4+}$ ions. These axes are defined so the $\hat{z}$ axis points along the local [111] direction and the $\hat{y}$ axis is oriented along one of the $C'_y$ axes perpendicular to the local [111], with frames on different basis sites related by $C_2$ rotations. If we consider rotations of the $d$ levels $U_r$ at each site $r$ that take the global cubic axes to the local frames then the operation $S$ acts in the local frame as $U_{S(r)}R_S U^r$. The set of quantization axes for the pyrochlore lattice has the advantage of acting only in the local frames, with the rotations of the local spin being the same across all the basis sites of the lattice up to a sign. Explicitly, one finds

$$P^{\hat{y}} U_{S(r)} R_S U^r \hat{y} = z_S, L_S,$$

where the operator $P$ projects into the $j_{\text{eff}} = 1/2$ subspace of the $d$ levels. The $z_{S,r}$ is a sign that only depends on the basis site of the pyrochlore lattice and can be found in Ref. [23] as the gauge transformations for the monopole flux state. The $L_S$ are spin rotations in the $\Gamma_{j=1/2} = \Gamma_{4g}$ representation of the site symmetry group $D_{4d}$ and be obtained from the generators

$$L_{C_3} = e^{-i\sigma^y z/3}, \quad L_{C_2} = L_4 = 1,$$

$$L_{C'_2} = L_{C'_4} = i\sigma^y$$

where $C_3$ and $C_2$ are independent of axis and the $C'_2$ and $C'_4$ are for the [110] and [100] axis respectively.

Here we will work only with the nearest neighbour hoppings, where aside from spin the hopping matrices depend only on the four basis sites. Extension to further neighbour hoppings is straightforward. Using symmetry operations in the local axes, $c^r \rightarrow z_{S,r} L_S c^{S(r)}$, one can show that there are only two allowed terms in the model

$$H_{Ir} = \sum_{\langle r'r \rangle} i c^r \left[ f_1 \gamma^+_r \gamma^+_r + f_2 (\sigma^x \gamma^+_r + \sigma^y \gamma^+_r) \right] c_{r'}.$$  

(4)

The $\gamma^+_r$ and $\gamma^+_r$ depend only the basis sites and can be written

$$\gamma^+_r = \begin{pmatrix} 0 & 1 & +\omega & +\omega \\ -1 & 0 & +\omega & -\omega \\ -\omega & -\omega & 0 & 1 \\ -\omega & +\omega & -1 & 0 \end{pmatrix},$$

$$\gamma^+_r = \begin{pmatrix} 0 & 1 & +1 & +1 \\ -1 & 0 & +1 & -1 \\ -1 & -1 & 0 & +1 \\ -1 & 1 & -1 & 0 \end{pmatrix},$$

where $\omega = e^{\pi i/3}$. Earlier studies have used formal global axes for the $j_{\text{eff}} = 1/2$ bands [24, 25], which can be obtained from the model derived above by inverting the local spin rotation $U_r$ at each site.

**Non-Kramers doublets and pseudo-spins:** Having established a model for the Ir$^{4+}$ ions, we now consider the Pr$^{3+}$ ions. Since these states are highly localized, being in a 4$f^2$ configuration, we use Hund’s rules to arrive at the ground state multiplet $3H_{4g}$, with inelastic neutron scattering studies of Pr$_2$Ir$_2$O$_7$ identifying a ground state doublet of $E_g$ character. The lowest lying excited state is a singlet $\sim 162k_{\text{B}}$ above the doublet, two orders of magnitude larger than the onset of the ordering, so we restrict to only the ground state doublet. This doublet has the form

$$|E_g, \pm \rangle = a_4 |\pm 4\rangle \pm a_1 |\pm 1\rangle - a_2 |\mp 2\rangle,$$

(5)

where $a_4$, $a_2$ and $a_1$ are real numbers depending on the details of the crystal field [27]. Within the space of doublets, super-exchange interactions are mediated through the surrounding oxygen atoms. This can be computed via a strong coupling expansion, including the effects of hopping between the Pr 4$f$ states and the O 2$p$ states. When projected into the subspace of doublets, the exchange Hamiltonian is most conveniently written using pseudo-spin operators

$$\tau^\alpha_r = \sum_{\alpha \beta} |E_{g, \alpha} \rangle \langle E_{g, \beta}| r \tau^\alpha \langle \alpha \beta|,$$

(6)

where $\alpha, \beta = \pm, \mu = x, y, z$ and $|E_{g, \pm}\rangle$ are the doublet states at site $r$. The $\tau^\alpha_r$ operator is magnetic, proportional to the magnetic dipole moment, while the transverse $\tau^x_r$ and $\tau^y_r$ parts are non-magnetic, carrying quadrupolar moments. All three exchanges allowed by symmetry are generated [28], giving the model in the local axes [27, 29, 30].

$$H_{Pr} = \sum_{\langle r'r \rangle} \left[ J_z \tau_z^r \tau_z^{r'} + \frac{J_z}{2} \left( \tau^\alpha_r \tau^\alpha_{r'} + \tau^\alpha_{r'} \tau^\alpha_r \right) \right]$$

$$+ J_{\pm} \sum_{\langle r'r \rangle} (\gamma^+_r \tau^+_r \tau^+_{r'} + \gamma^-_r \tau^-_r \tau^-_{r'}),$$

(7)
where $\tau^+_r = \tau^+_r \pm i \tau^0_r$ and the sums run over nearest neighbour bonds. Pseudo-spin rotational symmetry is not present when $J_z \neq J_{\perp}$, or in the presence of $J_{z\pm}$. The form of the $J_{z\pm}$ terms is a consequence of the intertwining of pseudo-spin and spatial symmetries, with the phases $\gamma_{rr'}$ defined as $\gamma_{rr'} = \hat{\gamma}_{rr'} \gamma_{rr'}^\dagger$.

Hybridization: We now consider interactions between the Pr and Ir sublattices, focusing on those mediated by hoppings between the sublattices, through physical or virtual processes. Charge transfer between the Pr and Ir necessarily involves intermediate states such as $4f^1$ or $4f^3$. For definiteness, we will assume that the $4f^1$ states are lower in energy than the $4f^3$, and thus dominate, though our results do not depend fundamentally on this choice. In the $D_{3d}$ crystal field this splits into a combination of $\Gamma_{4u}$, $\Gamma_{5u}$ and $\Gamma_{6u}$ representations. An example is the pair $\Gamma_{5u} + \Gamma_{6u}$, degenerate due to Kramers theorem, given by the $m = \pm 3/2$ states in $j = 5/2$ manifold of the $4f^1$ configuration. Hybridization between the the 5d $j_{\text{eff}} = 1/2$ states of the Ir and the localized states on the Pr can occur via several mechanisms, such as oxygen mediated hoppings, but an effective description written as direct hopping is possible once the intermediate states have been integrated out. Considering only intermediate states $\Gamma_{5u}$ and $\Gamma_{6u}$, the allowed hoppings are

$$H_{\text{hyb}} = V_z \sum_{rr'} \gamma_{rr'} e^{i \mathbf{r}_r \cdot \mathbf{r}_r'} / (1)_{r_r'} c^\dagger_{\mathbf{r}_r} [\Gamma_{5u}]_{r_r'} \langle E_g, \vec{\alpha} \rangle_{r_r'}$$

$$+ V_s \sum_{rr'} \gamma_{rr'} e^{i \mathbf{r}_r \cdot \mathbf{r}_r'} / (1)_{r_r'} c^\dagger_{\mathbf{r}_r} [\Gamma_{6u}]_{r_r'} \langle E_g, \vec{\alpha} \rangle_{r_r'} + \text{time reversed} + \text{h.c.}$$

where $\vec{\alpha} = -\alpha$, $r$ is an Ir site, $r'$ is a Pr site and $\gamma_{rr'} = \hat{\gamma}_{rr'} \gamma_{rr'}^\dagger$. If one splits the Ir-Pr bonds into two sets, related by inversion, then $(1)_{r_r'}$ is +1 on the first set and −1 on the second. This pattern is shown for the Pr centered hexagons in Fig. 2. To derive this form, one must keep in mind that the $\Gamma_{5u}$ and $\Gamma_{6u}$ states are Kramers and defined in the local axes, and so carry the same signs $\zeta_{S,r}$ as the $j_{\text{eff}} = 1/2$ states in their symmetry operations. For simplicity we set $V_s = 0$ for the remainder of this work, as it does not affect the results qualitatively.

We consider a fermionic slave-particle approach, as this allows for a natural treatment of hybridization between the Pr and Ir. The transition operators $[\Gamma_{5u}]_{r} \langle E_g, \vec{\alpha} \rangle_{r}$ and $[\Gamma_{6u}]_{r} \langle E_g, \vec{\alpha} \rangle_{r}$ are written using a pseudo-spinon $\eta_{r\alpha}$ and auxiliary bosons $\Phi_5$ and $\Phi_6$.

$$[\Gamma_{5u}]_{r} \langle E_g, \vec{\alpha} \rangle_{r} = \Phi_{5\alpha} \eta_{r\alpha},$$

$$[\Gamma_{6u}]_{r} \langle E_g, \vec{\alpha} \rangle_{r} = \Phi_{6\alpha} \eta_{r\alpha}$$

These slave-particles are constrained to satisfy $\eta_{r\alpha} \eta_{r\beta}^\dagger + \Phi_{5\alpha}^\dagger \Phi_{5\beta} + \Phi_{6\alpha}^\dagger \Phi_{6\beta} = 1$. Since these pseudo-spinons are of non-Kramers character, the symmetry operations in this local basis do not carry the signs $\zeta_{S,r}$ and transform simply as $\eta_{r} \rightarrow M_S \eta_{r}$ where $M_S$ is the pseudo-spin rotation corresponding to the symmetry operation $S$. The $\Phi_5$ and $\Phi_6$ bosons transform as the associated one-dimensional representations, but being Kramers states in the local quantization axes they also carry the phase factors $\zeta_{S,r}$ and transform as $\Phi_{5\alpha} \rightarrow \zeta_{S,r} e^{i \phi_{5\alpha}} \Phi_{5\alpha}$ and $\Phi_{6\alpha} \rightarrow \zeta_{S,r} e^{i \phi_{6\alpha}} \Phi_{6\alpha}$ under the symmetry operation $S$.

When splitting $\Delta$ between the $E_g$ and the excited states is large then we expect condensation of the bosons $\Phi_{5\alpha}$ and $\Phi_{6\alpha}$ at order $\Delta^{-1}$. In this limit the constraint can be simplified to $\eta_{r} \eta_{r}^\dagger \sim 1$. Condensing only in the $\Phi_5$ channel one has an effective hopping between electron $c^\dagger_\alpha$ and spinon $\eta_{r}$

$$H_{\text{hyb}} \sim V \sum_{rr'} \gamma_{rr'} e^{i \mathbf{r}_r \cdot \mathbf{r}_r'} / (1)_{r_r'} c^\dagger_{\mathbf{r}_r} \eta_{r\alpha} + \text{h.c.}$$

where we have absorbed $\Phi_{5\alpha}$ into $V_c$ defining $V = V_c \Phi_{s5}$. Having either the $\Phi_5$ or $\Phi_6$ channels to condense breaks time-reversal and time-reversal squared, an example of hastatic order.

Non-Kramers spin liquids: In terms of the slave-particles the pseudo-spin operator is given by $\gamma_{rr'} = \frac{1}{2} \eta_{r\alpha} \sigma^{\alpha\beta} \eta_{r'\beta}$. To render the problem tractable, we consider an approximate ground state generated from a Hamiltonian quadratic in the fermions.
Variational Monte Carlo calculations\textsuperscript{23, 33} on the Heisenberg model motivate us to consider two classes of \( U(1) \) spin liquid ansätze, the uniform and monopole states which are competitive in this limit. The monopol ansatz is a chiral spin liquid, breaking time-reversal and inversion but preserving the product, and can be characterized by hoppings carrying a flux of \( \pi/2 \) exiting the faces of each tetrahedron. The uniform state has equal hoppings on all bonds, carrying zero flux through all plaquettes. Since the presence of \( J_{\pm} \) or \( J_{+} \neq J_{-} \) breaks SU(2) pseudo-spin rotational symmetry, these ansätze must be extended using their respective projective symmetry group to include pseudo-spin-dependent \( E_{r}^{\text{hyb}} \) hoppings in addition to the pseudo-spin-independent \( \chi^{r}_{fr} \) hoppings allowed at the SU(2) symmetric point. Each spin liquid ansatz is characterized by a quadratic Hamiltonian

\[
H(\chi, E) = \sum_{\langle rr' \rangle} \left( \chi^{r}_{fr} \eta^{r}_{f} \eta^{r'}_{f} + \sum_{\alpha} E_{r}^{\alpha}_{f} \eta^{r}_{f} \sigma^{\alpha} \eta^{r'}_{f} \right),
\]

where the single occupancy constraint is implemented on average through chemical potentials \( \lambda_{r} \) tuned to enforce \( \langle \eta^{r}_{f} \eta^{r'}_{f} \rangle = 1 \).

To gain insight into which spin liquid may be favoured as we move away from the Heisenberg limit, for each ansatz Hamiltonian \( H(\chi, E) \) we compute the ground state \( |\psi(\chi, E)\rangle \). The energy \( \epsilon(\chi, E) = \langle \psi(\chi, E) | H_{irr} | \psi(\chi, E) \rangle \), where \( H_{irr} \) is the full Pr Hamiltonian, is then minimized with respect to \( \chi \) and \( E \). The phase diagram is shown in Fig. 3\textsuperscript{b} giving the state with lowest \( \epsilon \) as a function of \( J_{z}/J_{z} \) and \( J_{\pm} \). The monopol ansatz occupies large region of the phase diagram around the Heisenberg point, with \( E^{2} \) terms becoming finite at small \( J_{z} \) and the \( E^{z} \) components remaining disfavoured throughout. The uniform state is fully symmetric, with trivial PSG and does not become favoured until \( J_{\pm} \) is of order \( \sim J_{z}/2 \). The ansatz has the simple form \( \chi^{r}_{fr} = \chi, E_{r}^{\pm} = \gamma r r E \) and \( E_{r}^{r} = 0 \). We show the dispersion of this state when \( E \neq 0 \) and \( \chi = E/2 \) in Fig. 3\textsuperscript{a}. Note the lack of doubly degenerate bands, despite the presence of both time-reversal and inversion symmetry, due to these pseudo-spinons being non-Kramers.

**Broken Symmetries:** We now consider the full Hamiltonian

\[
H = H_{irr} + H_{hyb} + H_{P},
\]

adding in terms describing a uniform spin liquid on the Pr. When the \( \Phi_{r} \) boson condenses the \( U(1) \times U(1) \) gauge symmetry of the decoupled electron and spin-liquid system is broken to a single \( U(1) \)\textsuperscript{8, 34}, given by the transformation \( \eta \rightarrow e^{i\theta} \eta \) and \( c \rightarrow e^{i\theta} c \). This breaking of the relative gauge symmetry results in a Meissner-like effect, with a mass term pinning the emergent and physical gauge fields together. This pinning manifests in the acquisition of electric charge by the pseudo-spinon \( \eta \) and allowing the \( \eta \) pseudo-spinons to contribution directly to the Fermi sea as well as electromagnetic properties of the system\textsuperscript{55}. A more general problem, which can be accessed by considering further intermediate \( 4f^{3} \) and \( 4f^{3} \) channels, is an arbitrary hybridization

\[
H_{hyb} \sim \sum_{rr'} \sum_{\alpha \beta} V_{rr}^{\alpha \beta} \eta^{r}_{f} \eta^{r'}_{f} \eta^{r}_{f} \eta^{r'}_{f} + h.c
\]

Any choice of this \( V_{rr'} \), when both \( H_{irr} \) and \( H_{P} \) are present, will result in not only a breaking of time-reversal but in addition a breaking of at least one of the spatial symmetries. This is due to an incompatibility between the gauge structures of the Ir and Pr sublattices. For all operations \( S \) in \( Fd\tilde{S}m \), a symmetric hybridization must have

\[
V_{rr'} = \sum_{S \subset r} e^{i\alpha S} V_{S^{-1}(r)} S^{-1}(r) M_{S}^{r}
\]

for some choice of phases \( e^{i\alpha} \). Since the symmetries in the local axes form a group, for any operations \( S \) and \( S' \) the action of \( SS' \) must be equivalent to the action of \( S' \) followed by \( S \). The local rotations satisfy \( L_{S S'} = L_{S} L_{S'} \), so an equation relating \( z_{S S',r} \) to \( z_{S,r} \) and \( z_{S',r} \) can be obtained. Explicitly, this is given by

\[
z_{S S',r} = \eta_{S S'} z_{S,r} z_{S',r}^{-1}(r)
\]

with \( \eta_{S S'} = \pm 1 \). When combined with Eq. 13 this consistency condition entails that \( \eta_{S S'} \) be gauge equivalent to 1. For the PSG of \( z_{S,r} \) is false, and so this only satisfied by some subgroup of \( Fd\tilde{S}m \), breaking some of the symmetry.

The specific form shown in Eq. 1\textsuperscript{1}] motivated by the Anderson limit, breaks all spatial symmetries except for inversion and a single \( C_{3} \) axis. In the gauge used throughout the paper this is the [111] axis. As shown in Fig. 1\textsuperscript{5} if spin dependence is ignored, then we can understand the gauge structure in a qualitative fashion in a flux of \( \pi/2 \) exiting each tetrahedral face of the Ir sublattice. With the uniform spin liquid on the Pr and \( V_{rr} \) on the Pr-Ir bonds chosen as in Eq. 1\textsuperscript{1}] the flux is trapped in this truncated tetrahedron. Since the flux is not exiting, it must recombine into 2\( \pi \) flux somewhere within the volume. We have arranged it to preserve one of the \( C_{3} \) axes. When the Pr bonds are not present, this flux can cancel inside the remaining tetrahedra and thus form a symmetric state. In the presence of Pr-Ir bonds, a flux passes through the plaquettes between the Pr and Ir breaking the symmetries. With the only remaining symmetries are inversion and single \( C_{3} \) axis, the system is sufficiently asymmetric such that \( \sigma_{3} \) is allowed oriented along the [111] direction. Further magnetic, charge and quadrupolar orderings are generically induced, subject only to this fairly permissive \( C_{3} \) symmetry and inversion.

**Physical consequences:** To explore the effects of the spin liquid parameters and hybridization we fix \( t_{1} = 1, t_{2} = 0.1 t_{1} \) and \( \chi = E/2 \) and vary \( E \) and \( V \). This assumes that \( J_{\pm}/J_{z} \) is sufficiently large so that a uniform spin liquid is stabilized. Calculations of magnetization and AHE coefficients are shown in Fig. 1\textsuperscript{4} and Fig. 1\textsuperscript{5}. The magnetization shows the net magnetic moment per Pr atom, oriented along the [111], with contributions from both Pr and Ir sublattices (as shown in Fig. 1\textsuperscript{3} and 1\textsuperscript{4}) using \( g \) factors of \( g_{Pt} \sim 6.0 \) and \( g_{Ir} \sim 2.0 \). The anomalous Hall vector \( \Phi_{A} \) is computed using the Kubo
formalism[6], where the pseudo-spinons contribute as electrons to the current operators in the condensed phase.

The large AHE with small magnetic moments is in qualitative agreement with the properties of the intermediate phase of Pr$_2$Ir$_2$O$_7$. Here both the AHE and magnetic moment are considerably larger than observed experimentally. This discrepancy can be explained if the domains of the ordered phase are not fully aligned by the hysteresis process, then the observed AHE and magnetization would represent residual contributions from the partially aligned domains. At the mean field level one expects that the transition into the hybridized phase should show a jump in the specific heat. Since the order parameter can take 8 directions along the [111] axes one expects the critical theory to be described by an $O(3)$ type model, leading to a cusp at the transition. The effects disorder or Pr-Ir substitution can potentially smooth this cusp into the transition.

FIG. 4: (a-b) The AHE and net magnetization along the [111] direction for several values of $E$ as a function of $V$, with $t_2/t_1$ fixed at 0.1 and $\chi/E = 0.5$. (c-d) The pattern of local magnetic moments on the Ir and magnetic and quadrupolar moments on the Pr.

An essential feature of our proposal is the lack of large on-site moments. All induced orderings, such as the magnetic, quadrupolar and charge modulations are small, only appearing at fractions $\sim 10^{-2} - 10^{-3}$ of their saturated values. Distinct from scenarios with large moments that approximately cancel, leaving a small net moment. To distinguish these experimentally, nuclear magnetic resonance (NMR) on the oxygens is promising. In the crystal structure the oxygens lie in two inequivalent Wyckoff positions: the 8$a$ position, which has tetrahedral symmetry and the 48$f$ position which is in an asymmetric location. If one can account for small net moment through a cancellation of large local moments, then one may expect the net field at the symmetric site 8$a$ to be small, but generically asymmetric site 48$f$ should be affected by a net field from moments of order $\sim \mu_B$. In this scenario assuming dipolar fields from the moments acting at the oxygen sites, one then expects the effect at the 48$f$ site to be several orders of magnitude larger than that of the 8$a$ site. Our proposal predicts a significantly different result, with the small local moments producing only small fields of order $\sim 0.1 - 1G$ at both oxygen sites.

Conclusion: We have proposed a mechanism for symmetry breaking when conduction electrons hybridize with a quantum spin liquid. Applied to Pr$_2$Ir$_2$O$_7$ we found that the hybridization of two subsystems ($f$ and $d$ electrons) results in a chiral nematic metal with broken time reversal and spatial symmetries, exhibiting an anomalous Hall effect without a sizeable magnetic moment. This mechanism could potentially manifest in a wide range of heavy fermion materials on geometrically frustrated lattices.

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