Spinon confinement and a sharp longitudinal mode in Yb$_2$Pt$_2$Pb in magnetic fields

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The fundamental excitations in an antiferromagnetic chain of spins-1/2 are spinons, de-confined fractional quasiparticles that when combined in pairs, form a triplet excitation continuum. In an Ising-like spin chain the continuum is gapped and the ground state is Néel ordered. Here, we report high resolution neutron scattering experiments, which reveal how a magnetic field closes this gap and drives the spin chains in Yb$_2$Pt$_2$Pb to a critical, disordered Luttinger-liquid state. In Yb$_2$Pt$_2$Pb the effective spins-1/2 describe the dynamics of large, Ising-like Yb magnetic moments, ensuring that the measured excitations are exclusively longitudinal, which we find to be well described by time-dependent density matrix renormalization group calculations. The inter-chain coupling leads to the confinement of spinons, a condensed matter analog of quark confinement in quantum chromodynamics. Insensitive to transverse fluctuations, our measurements show how a gapless, dispersive longitudinal mode arises from confinement and evolves with magnetic order.
The one-dimensional (1D) XXZ Hamiltonian for quantum spin chains given by Eq. (1) is a cradle of exactly solvable quantum theory models of interacting many-body systems\(^6\). The Hamiltonian considers the components \(S_i^x = x, y, z\) of a spin angular momentum operator, \(S_i^z = 1/2\) at site \(i\) on a 1D chain, with \(J\) the nearest neighbor exchange coupling for \(x, y, z\) spin components, \(\Delta\) a uniaxial coupling anisotropy, and \(H\) magnetic field (with \(g\) and \(\mu_B\) the Lande g-factor and Bohr magneton respectively),

\[
\mathcal{H} = J \sum_i \left( S_i^x S_{i+1}^x + S_i^y S_{i+1}^y \right) + \Delta S_i^z S_{i+1}^z - g\mu_B \sum_i H \cdot S_i. \tag{1}
\]

The low energy excitations of this model Eq. (1) are spin-1/2 quasiparticles called spinons. In the limit of strong Ising anisotropy, \(\Delta \gg 1\), spinons can be visualized as domain walls in an antiferromagnetically ordered ground state of the chain (Fig. 1a). Angular momentum conservation mandates that spinons are always created in pairs, such that each spinon carries a fraction, \(\pm 1/2\), of the angular momentum change, \(\Delta S^z = 0\), \(\pm 1\), required to initially introduce the domain walls in an infinite chain. Since moving these domain walls is an energy and angular momentum conserving process, the walls will propagate freely, carrying the quanta of energy, \(q_0\), and linear momentum, \(q\), introduced by their creation (Fig. 1a). The physics contained in Eq. (1) leads directly to the separation of the spin from other electronic degrees of freedom, mapping directly onto that of the Luttinger liquid for \(-1 \leq \Delta \leq 1\).

Coupling the chains described by Eq. (1) leads to new and emergent physics. Analogous to quark confinement in quantum chromodynamics\(^5,9\), the dimensional crossover from 1D chains to 3D coupled chains leads to quasiparticle confinement, thereby stabilizing long range magnetic order at temperature \(T > 0\). A new excitation of the longitudinal degree of freedom of the order parameter is predicted when the interchain coupling is weak\(^9\).

These phenomena have been the subject of a considerable amount of recent experimental work in XXZ spin chain materials\(^11\). Like a similar longitudinal mode previously observed near the critical point in a system of coupled spin-1/2 dimers\(^16,17\), this excitation can be interpreted as a condensed matter analog of the Higgs boson\(^18\).

Here, we report neutron scattering experiments on the one dimensional rare-earth antiferromagnet Yb\(_2\)Pt\(_2\)Pb to investigate these fundamental processes in detail, using an external magnetic field as a tuning parameter.

**Results**

**Inelastic neutron scattering on Yb\(_2\)Pt\(_2\)Pb.** Yb\(_2\)Pt\(_2\)Pb is a metal with a planar crystal structure where orthogonal pairs

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Fig. 1 Spinons on decoupled chains. **a** An anisotropic AFM spin chain (top). If two spins are interchanged, two domain walls are created between the original antiferromagnetic domain (green) and a new domain (blue) (middle). Those domain walls form the basis for the propagating states carrying the spinon gap \(E_0\), summed over \(0 \leq q_{\text{flu}} \leq 1\) rlu. The lower boundary of the spectrum (white circles) is shown along with the boundaries of the two spinon continuum obtained by fitting the lower boundary (red lines, \(\Delta = 3.46\)) and comparing the total measured spectrum to theory (white lines, \(\Delta = 2.6\))\(^26\). The color scale for parts **b**–**d** is shown above part **c**. Error bars represent one standard deviation. **b** The magnetic excitation spectrum and dispersion along the \(q_{\text{flu}}\) direction in reciprocal space measured at \(T = 0.1\) K, summed over \(-1 \leq q_{\text{flu}} \leq 1\) rlu. **c** The dispersion of particle-like (red) and hole-like (black) spinons, symmetric about \(E = 0\) in zero magnetic field, are sketched with the real \(\Delta = 2.6\) parameters for Yb\(_2\)Pt\(_2\)Pb. The bandwidth parameter \(I\) and the spinon gap \(\Delta_s\) are indicated by arrows, with \(2\Delta_s\) the energy separation between the particle and hole bands at \(q_{\text{flu}} = 0, 1\), and 2 rlu.
of Yb ions form a Shastry-Sutherland lattice (SSL) motif in the tetragonal a-b plane. High resolution neutron scattering experiments recently showed that the physics of 4f-orbital overlaps leads to unusual consequences for the magnetism in Yb$_2$Pt$_2$Pb. The low energy magnetic excitations are spinons, having a quantum continuum for momentum along the chain direction, which can be measured with inelastic neutron scattering, with an excitation bandwidth that is considerably larger than the excitation gap (Fig. 1b). At zero magnetic field, our measurements agree well with time-dependent density matrix renormalization group (tDMRG) calculations for the XXZ model Eq. (1) (Fig. 1d), although experiment indicates that the spectral weight is spread throughout the spin Brillouin zone (BZ) more evenly and to higher energies than these calculations predict, suggesting non-negligible next-neighbor coupling. Comparisons of our data to theory indicate only a modest anisotropy, $J_1 > J_2$. It is clear that the XXZ Hamiltonian Eq. (1) is an appropriate description for Yb$_2$Pt$_2$Pb despite the large and orbitally dominated moment of the Yb ions. Due to their Kramers doublet ground state of almost pure $|J, m_J\rangle = \{7/2, \pm 7/2\}$, the Yb moments have a pseudospin $S_z = 1/2$ character, $\Delta\neq 0$. Rather than quenching the quantum spin dynamics, the strong Ising magnetic anisotropy imposed by the crystal electric field acting on the $f$-orbital wave function instead singles out the longitudinal excitation channel in two orthogonal sublattices of 1D chains with momenta oriented along the (110) and (110) crystal directions.

The essential features of the quantum continuum can be understood by noting that each spinon carries spin 1/2, and so the angular momentum selection rules dictate that neutron scattering in Yb$_2$Pt$_2$Pb measures two-spinon states where the total spin is zero, with one spinon in each spin state, $\pm 1/2$. In order to describe the boundaries of the two-spinon continuum, it is convenient to adopt the language of particles and holes occupying the fermionic spinon dispersion along the chain direction, $E_{ph} = \pm \left(\hbar^2 \sin^2(\pi q_{\text{fl}}) + \Delta_0^2 \cos^2(\pi q_{\text{fl}})\right)^{1/2}$, $0 \leq q_{\text{fl}} < 1$ rlu, where $\Delta_0$ is an energy gap brought on by the XXZ anisotropy $\Delta > 0$, and $I$ defines the dispersion bandwidth and encodes the coupling $J$ (Fig. 1e). In place of electric charge, these particles and holes each carry a half unit of spin angular momentum. The boundaries of the two-spinon continuum are defined by the extremal energy and momentum conserving combinations of one particle and one hole, and they are shown in Fig. 1b for both $\Delta = 2.6$ and 3.46, the range of values determined in previous work. (see Supplementary Note 1).

At zero magnetic field, the chemical potential is in the middle of the gap separating the particle and the hole bands, which describes the antiferromagnetic (AFM) state with zero total spin, $S_z = 0$. The size of the $T = 0$ ordered moment implied by the XXZ anisotropy is consistent with our measurements at $T = 0.1$ K, within the precision of our data. This implies that the interchain coupling responsible for moving the Néel temperature away from $T = 0$, the value predicted by the XXZ model, to $T_N = 2.07$ K is less than both the intrachain exchange $J = 0.206$ meV and the spinon gap $\Delta_0 = 0.095$ meV, the dominant 1D energy scales. The flat dispersion of the excitations between the chains in zero field (Fig. 1c), despite the apparent ladder geometry of the crystal structure, suggests that the effect of interchain interactions on low energy excitations is quenched when $\Delta_0$ is nonzero.

A magnetic field along the z (110) direction introduces the Zeeman term $-g_\mu_B H \sum_i S_z^i$ to Eq. (1), which lowers the chemical potential, $\mu = -g_\mu_B H S_z^2$. The potential needed to close the energy gap for creating a hole on the spinon dispersion is $|\mu| = \Delta_0 = 0.095$ meV. Taking $g = 7.3$, $|\mu| = \Delta_0$ corresponds to a critical field of $\mu_B H = 0.5$ T. An abrupt increase in the bulk magnetization is seen at this field when oriented parallel to the magnetic moments of either sublattice at temperatures $k_B T < \Delta_0$ (Fig. 2a)\textsuperscript{23,24}. On the other extreme of the magnetization, when $|\mu| > I$, the entire hole band lies above the chemical potential, the field having transformed all holes to particles. Particle-hole pairs can no longer be produced, quenching spinon excitations, and producing a ferromagnetic (FM) state. The saturation field in Yb$_2$Pt$_2$Pb is 2.3 T, precisely the field needed for $\mu = 0.485$ meV, the value of $I$ when $\Delta = 2.6$, the number obtained by requiring that Eq. (1) provides best description of the entire $\mu_B H = 0$ excitation spectrum\textsuperscript{25}. The comparison is less favorable when $\Delta$ is taken to be 3.46, the value derived directly from fitting the lower boundary of the continuum. As the static properties correspond to an integration over all energies, it is not surprising that they are better captured by $\Delta = 2.6$ and we adopt this value of $\Delta$ here.

At intermediate fields, $0.5 < \mu_B H < 2.3$ T, the hole band is partially emptied. The chemical potential crosses the hole dispersion at four points in the spinon BZ (Fig. 2b), defining a Fermi wavevector $k_F$ that directly links particle and hole states. There are now eleven unique extremal states made from a single particle and hole, rather than the three that are possible in zero field. The boundaries of the two spinon continuum change dramatically, as the possible extremal states are heavily influenced by the restrictions of the hole energies and momenta and the additional phase space occupied by particles.

That is precisely what is measured in the neutron scattering spectra of Yb$_2$Pt$_2$Pb (Fig. 2d-f). At $\mu_B H = 1.0$ T, there is strong scattering concentrated at low energies within a range $\pm 2k_F$ around the BZ center, with a weaker continuum at higher energies. As the magnetic field is increased, the low-energy spectral weight spreads throughout the zone as $k_F$ increases, with higher energy pockets of spectral weight bounded by the extremal two spinon states comprising the continuum. While the measured continua are in broad agreement with the results of tDMRG calculations performed on isolated chains (Fig. 2g-i), there are marked differences at low energies where interchain interactions are important. The measured lower boundaries are gapped at small energies and also slightly distorted relative to theoretical expectations, with the increased spectral weight indicative of a bound state. This is a direct demonstration of spinon confinement induced by interchain coupling, which is not accounted for in the 1D calculations of Fig. 2g-i, and also shows how a magnetic field tunes the coexistence of confined and free spinons in Yb$_2$Pt$_2$Pb for $\mu_B H > 0.5$ T.

The schematic picture of spinons and their propagation presented in Fig. 1a needs to be modified in the presence of the interchain interactions, where the creation of spinons on one chain leads to frustration of the AFM interactions between the chains. As two spinons separate, the energy of this frustration grows with the number of FM aligned neighbors (Fig. 2c). This provides a linear confining potential, just as quarks are confined by the gluon-mediated strong force in QCD, which also increases with quark separation. When spinons are created with energies above the highest energy level existing in the confining potential introduced by the interchain coupling, these high energy quasiparticles propagate freely within the two spinon continuum, demonstrating the same asymptotic freedom as experienced by unbound quarks. A spinon bound state is observed in a neutron scattering experiment as excess spectral weight of resolution-limited energy width, which prominently appears near $\pm 2k_F$, the two soft spots around the BZ center, below the quantum continuum.\textsuperscript{10}
Perhaps the most interesting aspect of these free and confined excitations is how their dispersions develop in the $q_{HH}$ direction, perpendicular to the chains. Spinons are continually created in registry on adjacent chains, thus minimizing the inter-chain dispersion. Most notably, for $0.5 \lesssim \mu_0H \lesssim 2.3$ T, we observe a new excitation that emerges from the featureless spinon continuum found along $q_{HH}$ in the gapped zero field Néel phase where the chains are effectively decoupled. This mode resides within the low-energy window of the two spinon bound states, but has a pronounced dispersion in the $q_{HH}$ (interchain) direction that changes considerably with increasing field (Fig. 3b–d). Remarkably, at 1.0 T the dispersive interchain mode appears nearly gapless, while its intensity is markedly larger than that of the continuum, Figs. 2d and 3b. The mode becomes clearly gapped with increasing fields, while the

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**Fig. 2** Spinons in a magnetic field. a The magnetic field dependence of the static magnetization at $k_B T < \Delta$, [red, right axis, $H_{\parallel(110)}$] shows several discontinuous jumps (blue and green dashes), corresponding to transitions among different 3D ordered phases that are more clear in $dM/d\mu_0H$ (black, left axis)\(^2\). These phases correspond to different ways that magnetic moments arrange into registry minimizing the energy of magnetic dipole interactions between the Yb moments. b When a magnetic field is applied along the chain direction, the chemical potential $\mu = -\mu_0H S^z$ (yellow) is lowered, emptying part of the hole band when $|\mu| > |\Delta|$. $\mu$ crosses the hole dispersion at four points in the Brillouin zone (black arrows), defining the Fermi wavevector $k_F$. c Two AFM ordered, 1D spin chains (top). If two spins on one chain are interchanged, two domain walls are formed between the original domain (green) and a new one (blue). The new domain frustrates the interchain interaction, represented by the red interchain bonds. This frustration creates a linear potential confining low energy spinons to bound states (bottom). d-f The magnetic excitation spectrum and its dispersion along the $q_{L}$ direction in reciprocal space measured at $T = 0.1$ K and $\mu_0H = 1.0$ T (d), 1.5 T (e), and 1.7 T (f), summed over $-1 \leq q_{xy} \leq 1$ rlu. The dispersions for the extremal combinations of particles and holes are shown (black lines) (See Supplementary Note 2). The spinon bound states are manifest from the enhanced low energy spectral weight around $1 \pm 2k_F$ (black arrows). g-i The spinon spectrum computed using 1DMRG calculations for the XXZ model Eq. (1) on a 96-site chain with $\Delta = 2.6$ at equivalent chain magnetizations as (d-f), shown on the same color scale. The dispersions for the extremal combinations of particles and holes are also shown as black lines.
relative spectral weight of the continuum grows. We model the energy dependence of the scattering at a specific \(q_{\text{HH}}\) as a damped harmonic oscillator (DHO) response centered at the mode position and the product of a Lorentzian and step function accounting for the continuum at higher energies, all convolved with the instrument resolution (Fig. 3e). The energy width of the new mode is roughly resolution limited at all fields, and it is always distinguishable from the spinon continuum for fields \(\mu_0 H \leq 1.7 \text{T}\). The connection of the mode and the confined spinon states can be emphasized by integrating over the energies of the mode and plotting the intensity as a function of momentum along the chains (Fig. 3e-Inset). At all fields, \(\approx 85\%\) of the...
spectral weight is concentrated within the momentum range \( q_L = 1 \pm 2k_F \).

Importantly, this interchain mode is longitudinally polarized. We confirm its longitudinal character by using the fact that neutron scattering cross-section is uniquely sensitive to magnetic fluctuations that are perpendicular to the wave vector transfer\(^{27}\). The intensity measured at 4 T, in the FM state precisely follows the projection of the scattering wave vector on the (110) direction, revealing fluctuations polarized along the in-plane Ising moments, which are insensitive to magnetic fields (Fig. 3). When this field-independent contribution is subtracted as a background, the resulting field-dependent intensity \( I_{\mu H}(q) \) does not depend on the wave vector orientation in the scattering plane, indicating magnetic fluctuations polarized along the vertical direction, collinear with the magnetic field\(^{26}\) (see Supplementary Note 5). At all fields, our measurements unambiguously probe the longitudinal response. The Ising anisotropy of the 7/2, \( \pm 7/2 \) ground state doublet of the Yb moments nearly completely suppresses any transverse magnetic fluctuations from our measurements.

The longitudinal interchain mode changes dramatically over a relatively narrow range of fields as the underlying antiferromagnetic order is weakened and ultimately destroyed. The low temperature \( H - T \) magnetic phase diagram of Yb\(_2\)Pt\(_2\)Pb (Fig. 3a) has several different AFM ordered phases\(^{22-25}\). In zero field, there is a five by five periodicity to the order in the tetragonal a – b plane, evidenced by neutron diffraction peaks at index as \( q_{HH} = 0.2 \text{ rlu} \) (Fig. 4a)\(^{29}\). When the gap \( \Delta_S \) closes at \( \mu_0H = 0.5 \text{ T} \), those peaks move from \( q_L = 1 \text{ rlu} \) to incommensurate positions along \( q_L \) (Fig. 4b), consistent with the longitudinal component of the spin-spin correlation function probed by our neutron scattering measurements being locked to twice the Fermi wave vector\(^{10}\). The ordering wave vector follows 2\( k_F \) in turn, connecting to the softest parts of the excitation spectrum.

![Fig. 4](image_url) Magnetic order and phase diagram of Yb\(_2\)Pt\(_2\)Pb. a – d. Time of flight data with energy transfer \( E = 0 \) in the \( q_{HH} - q_L \) plane, measured at \( T = 0.1 \text{ K} \) in fields of 0.025 T (a), 0.75 T (b), 1.2 T (c), and 1.7 T (d). The magnetic Bragg diffraction peaks in each part of the phase diagram are labeled 1-4. e, f The location in reciprocal space along \( q_L \) (e) and \( q_{HH} \) (f) of the magnetic Bragg scattering as a function of the field in the regions plotted in parts a – d. Data shown are an average, symmetrized about \( q_L = 1 \) and \( q_{HH} = 0 \). The open symbols correspond to the peaks labeled 1-4 in panels b – e, with black indicating peak 1, blue peak 2, red peak 3, and green peak 4. Magnetization divided by the saturation magnetization \( M/M_{sat} \) as a function of the field along the (110) crystal direction is also shown in panel e (pink line, right axis), demonstrating the initial trend, \( M/M_{sat} \approx (q_L - 1) \) and the concurrence of the jumps in magnetization with the abrupt changes in the position of elastic scattering. g The intensity of the Bragg peaks in parts b – f. The AFM order that emerges for \( \mu_0H > 1 \text{ T} \) is considerably weaker than the low field order, while the low field order falls off very abruptly for \( \mu_0H > 1.2 \text{ T} \). Symbols are the same as in panels e, f. h The zero field magnetic structure and dipole interactions in Yb\(_2\)Pt\(_2\)Pb. Yb moments in five unit cells are shown along the diagonal (110) direction, with the interchain couplings, \( J_{\perp}^{zz} \), as indicated. The Yb SSL AFM layers are highlighted (blue and green arrows), with the periodicity given by FM pairs every five unit cells (black and red arrows).
There are several small and abrupt shifts in the ordering wave vector along both $q_i$ and $q_{fi}$ (Fig. 4a–f) that coincide with abrupt jumps in the derivative of the low temperature magnetization [Figs. 2a and 4c, f], which manifest changes in 3D magnetic ordering as the magnetic moments re-arrange to minimize the energy of magnetic dipole interactions. For fields $\mu_0H > 1.0$ T, there is an emergence of a second incommensurate AFM ordered phase that is accompanied by the swift collapse of the original five by five order for $\mu_0H > 1.2$ T and even a third incommensurate order that persists up to the saturation field (Fig. 4c–g). As the new longitudinal interchain mode is an excitation of the underlying order, it is not surprising that it changes so dramatically between 1.0 and 1.5 T.

For $0.5 \leq \mu_0H \leq 1$ T, the $q_i$ component of the magnetic Bragg peaks follows the magnetization (Fig. 4e), which reflects the position of the spinon Fermi wavevector, $2k_F$ (Fig. 2b) (See Supplementary Note 6). Theoretical description of the excitation spectrum in this spin-density-wave (SDW) phase can be obtained by applying bosonization methods for quasi-1D spin-1/2 antiferromagnets. The low energy sector of such an antiferro- order terms. These changes are consistent with the observation from the magnetic diffraction (Fig. 4) of a tendency towards a weaker, frustrated longitudinal antiferromagnetic order. Smaller changes are found in the higher order terms. These changes are consistent with the observation from the magnetic diffraction (Fig. 4) of a tendency towards a weaker, frustrated longitudinal antiferromagnetic order as the applied magnetic field progressively polarizes the moments (See Supplementary Note 7).

### Discussion

It is difficult to visualize the nature of the interchain mode itself, as it is purely quantum mechanical in its origin, with no simple classical analog. In a conventional 3D ordered magnet, these interchain excitations would be transverse spin waves—pseudo-Goldstone modes of an antiferromagnetic order parameter, acquiring a small gap (mass) in the presence of spin anisotropy. The longitudinal polarization reveals that the new excitations observed here in Yb$_2$Pt$_2$Pb, which are separated from $E = 0$ with a field-dependent gap $\delta < 0.12$ meV, are in fact far more exotic. They represent amplitude excitations of the AFM order parameter, i.e., the staggered magnetization, analogous to the amplitude modes of the superconducting order parameter found in NbSe$_3$, and the Higgs boson. There is a large literature on the subject in the context of the theory of quantum magnets (e.g., refs. 33, 34, 37, 38), but so far there have been few experiments among weakly coupled chain systems that probe this mode and its dispersion across different regimes of interchain coupling, or its decay into transverse magnons in detail. This damping causes the longitudinal mode to appear as a resonance in the longitudinally polarized continuum rather than the dispersing quasiparticle like excitation observed here in Yb$_2$Pt$_2$Pb, which often obscures the physics entirely and has even led some to question the assumptions behind the theory. Recently, some evidence for a sharp longitudinal mode coexisting with the transverse spin waves has been obtained via polarized neutron scattering measurements in a more strongly coupled 2D ladder system, while other materials with similar XXZ anisotropy to Yb$_2$Pt$_2$Pb tend to confine all spinons into many modes. Here we overcome the limitations of such experiments thanks to the tuning parameters of Yb$_2$Pt$_2$Pb, which uniquely single out the longitudinal channel and allow us to clearly identify the dispersing amplitude mode and the deconfined 1D excitations at higher energies.

The physics of Yb$_2$Pt$_2$Pb that suppresses transverse spin waves also protects these longitudinal excitations and allows detailed observation of this mode dispersion and its dependence on an applied magnetic field, which tunes the ordered state across different phases. Accompanying transverse excitations—the spin
waves—must exist, but are not observed (Fig. 3f). They are suppressed by a factor \((q_i^0)^2 \gtrsim 100\) and cannot be measured in this scattering geometry, where even the application of a substantial magnetic field leaves behind the longitudinal continuum from magnetic moments perpendicular to the field. While further experiments are needed in alternate scattering geometries to clearly observe the accompanying transverse excitations, the present results quantify in unique detail an unusual dispersing longitudinal mode, a Higgs-like excitation of a effective spin-1/2 longitudinal mode, a Higgs-like excitation of an effective spin-1/2 magnetic field.

**Methods**

**Neutron scattering.** The \(\text{Yb}_2\text{Pt}_2\text{Pb}\) sample used in these measurements was the same sample used in ref. 26. The sample consists of approximately 400 co-aligned \(\text{Yb}_2\text{Pt}_2\text{Pb}\) single crystals (total mass \(\approx 6\) g) mounted to aluminum plates. For all measurements the sample was oriented with the (1 1 1) crystal direction vertical leaving the (110) and (001) directions in the horizontal scattering plane, making the scattering plane (\(H, H, L\)) in reciprocal space. All momenta are given in reciprocal lattice units (rlu), with 1 rlu given by \(2\pi a/\lambda\), where \(a = 5.776\ \text{Å}\) along the \(H\) direction and \(2\pi c/\lambda = 2\pi n/2.02 \approx 0.895\ \text{Å}^{-1}\) along \(L\). The crystallographic unit cell is twice the Yb-Yb near neighbor spacing along the \(c\)-axis—the relevant spacing for spinons. Therefore, the Brillouin zone for spinons is indexed from 0 to 1 rlu along (0, 0, 1), rather than the typical 0 to 1 rlu. Notationally, \(q_{\text{lt}}\) is parallel to the (\(H, H, 0\)) direction, with scattering primarily coming from \((H, H, 1)\) while \(q_{\text{lt}}\) is along the (0, 0, 1) direction.

The inelastic neutron scattering measurements on \(\text{Yb}_2\text{Pt}_2\text{Pb}\) making up the bulk of the data in this paper were made on the OSIRIS spectrometer at the ISIS neutron source at Rutherford Appleton Laboratory in Didcot, Oxfordshire, UK. The sample was mounted in a dilution refrigerator inside of a 7 T field of 0.025 T was used to suppress superconductivity in the aluminum sample holder. In general, inelastic neutron scattering probes the dynamical spin correlation function. Because of the crystal field ground state doublet of the Yb ions in \(\text{Yb}_2\text{Pt}_2\text{Pb}\) we are sensitive only to the longitudinal component of this function (see Supplementary Note 5).

Because of the detector coverage does not include an entire Brillouin zone, \((0 < q_{\text{lt}} < 2\text{ rlu}, 0 < q_{\text{HH}} < 1\text{ rlu})\) we integrate all data in the scattering plane in the horizontal plane unaffected due to the ground state doublet of the Yb ions. Measurements made at \(\mu_0 H = 4\) T can therefore be used as a background for measurements made at \(\mu_0 H = 2\) T, isolating only the lower field contribution from magnetic moments oriented parallel to the field. All neutron scattering results from OSIRIS have a measurement made at \(T = 0.140\) K and \(\mu_0 H = 4\) T subtracted in this fashion. For the nominal zero field measurements, the same sample was mounted in the same scattering geometry as the OSIRIS spectrometer at the ISIS neutron source at the SISIS neutron source at https://doi.org/10.5286/ISIS.E.42580328.

The data that support the results of this study are available from the corresponding authors upon reasonable request. Original time-of-flight data is available for the experiments at theISIS neutron source at https://doi.org/10.5286/ISIS.E.42560328.

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