The series of rare earth (R) ternary silicides, R$_2$PdSi$_3$, is characterized by multiple competing interactions, which give rise to complex magnetic and electrical behavior. Members of this series order antiferromagnetically with Néel temperatures between 8 and 21 K [1], however some members display different behavior. For instance, Nd$_2$PdSi$_3$ displays anomalous ferromagnetism, explained by the hybridization of the local 4$f$ moments with the conduction electrons [2], and, notably, Gd$_2$PdSi$_3$ displays numerous poorly understood anomalies, such as a resistivity minimum at 45 K which is suppressed by magnetic fields, negative magnetoresistance [3], which suggests a complex Fermi surface, non-linear Hall [4] and Nerst effects [5], and multiple step-like metamagnetic transitions at various fields below the zero-field ordering temperature, $T_N = 22$ K [6]. Interestingly, these properties display highly anisotropic behaviour, despite Gd$^{3+}$ being an s-state ion.

In a series of recent reports by Kurumaji et al. [4, 5, 7], the authors suggest that in Gd$_2$PdSi$_3$ at temperatures below 15 K and magnetic fields between 0.5 and 1.2 T parallel to the c axis a hexagonal lattice of magnetic skyrmions [A phase in Fig. 1(a)] is stabilized, flanked by two other poorly understood incommensurate phases termed IC-I at lower fields and IC-II at higher applied fields. Magnetic skyrmions are non-coplanar topological whirlpool-like spin structures with many proposed spintronic applications from data storage and logic devices to stochastic and neuromorphic computing [8–10].

Conventionally, the stabilization of magnetic skyrmions requires the Dzyaloshinskii-Moriya interaction [8, 11–18]. However, this interaction vanishes in Gd$_2$PdSi$_3$ due to the centrosymmetry of the hexagonal AlB$_2$-type crystal structure, shown in Fig. 1(b). In this structure, the magnetic Gd$^{3+}$ ions are organized within triangular planes perpendicular to the c direction, intercalated between a honeycomb arrangement of Pd and Si$^{2-}$ ions that order in a superstructure along the c axis. [19]. This structure gives rise to geometric frustration of the magnetic ions that, together with the Ruderman-Kittel-Kasuya-Yosida (RKKY) and higher-order biquadratic exchange interactions, could potentially stabilize both isolated skyrmions and skyrmion lattices, as well as a further vast variety of other other multi-q, noncollinear magnetic orderings [20–23].

Ultimately, the origin of the stability of the magnetic spin textures in Gd$_2$PdSi$_3$ is not clear. An angle-resolved photoemission spectroscopy study revealed Fermi surface nesting along the magnetic wavevectors, and suggested that a nesting-driven enhancement of the RKKY interaction stabilizes the incommensurate magnetic structures [24]. However, a second study study using local force calculations alternatively suggested that inter-band frustration stabilizes the zero field states, rather than the RKKY interaction [25], consistent with recent anomalous magnetoresistance results [3].

In other systems with well-nested regions in the Fermi surface, the itinerant electrons typically become unstable towards the formation of a charge and/or spin density wave (CDW/SDW) by a variety of mechanisms [26]. For instance, in elemental chromium the presence of near-perfect nesting conditions couples the itinerant electrons to phonons located at opposite regions of the Fermi surface, forming a SDW with twice the nesting vector [27]. Another possible stabilisation mechanism for a CDW state is the coupling of electrons to an incommensurate ordering of local moments inducing a S/CDW with the same period as the local-moment magnetic order [28]. These mechanisms are not exclusive; a cooperative interaction was found in GdSi where the local moments provide the necessary coupling between the nested regions of the Fermi surface to induce a SDW in the itinerant electrons, meanwhile the maximum in the itinerant ele-
tron susceptibility simultaneously induces local moment ordering with the same wavevector [29]. The possible appearance of such CDW states in Gd$_2$PdSi$_3$ has not previously been rigorously considered.

In this Letter, we investigate the incommensurate textures within Gd$_2$PdSi$_3$ using REXS and AC magnetometry. We demonstrate the presence of an incommensurate CDW state and coplanar magnetic textures at low temperature in magnetic fields up to 1 T. These magnetic textures are not compatible with the previous suggested helical and skyrmion textures, and we discuss potential coplanar alternatives. We propose that the observed anomalously-large nonlinear Hall effect in the A-phase is not due to the skyrmion topological Hall effect (THE) and must have a different origin.

A $2 \times 2 \times 5$ mm$^3$ single crystal of Gd$_2$PdSi$_3$ was grown by the optical floating zone technique at the University of Warwick. X-ray Laue diffraction determined the sample orientation and a high level of crystallinity was observed. Magnetic susceptibility with magnetic field $H \parallel c$ and a drive field of 1 mT at 10 Hz was measured using a Quantum Design MPMS3 at the I10 support laboratory, Diamond Light Source. The phase diagram was mapped out using scans of increasing field from 0 to 2 T after zero field cooling (ZFC) the sample from 30 K at each temperature. REXS from three satellite peaks surrounding the crystalline (300)-Bragg peak at the Gd L$_2$ edge (7.931 keV) was performed at the I16 beamline of Diamond Light Source, with incoming x-rays linearly polarized within the scattering plane [$\pi$ polarization on Fig. 1(d)]. Polarization of the outgoing scattered x-rays was determined using a gold analyzer crystal to separately measure the outgoing scattering polarised either within the scattering plane, $\pi$-$\sigma$, or perpendicular to the plane, $\pi$-$\pi$. In our experiment, the x-ray out-of-plane ($\sigma$) direction was collinear with the crystallographic c-axis. Field scans were taken after a ZFC protocol, up to the calibrated maximum field of 0.95 T. Temperature control was achieved using a $^4$He cold finger cryostat with helium exchange gas inside a beryllium dome. A magnetic field was applied using a water-cooled copper coil electromagnet with iron pole pieces.

The $H$-$T$ magnetic phase diagram (H $\parallel$ c axis) for Gd$_2$PdSi$_3$ measured using the real component of AC-susceptibility is shown in Fig. 1(a), and is consistent with previous reports. Below the magnetic ordering temperature of 21 K, we find clear maxima in the susceptibility delineating the phase boundaries between the three magnetically ordered phases, labelled IC-I/II (incommensurate) and A (previously suggested skyrmions). We note that the A phase is not characterized by the dip in the susceptibility that would commonly be associated with the skyrmion state in other materials [30]. These magnetic phases have been shown to exhibit incommensurate components, described by magnetic wavevectors $\mathbf{\bar{q}}_n$, which propagate along the reciprocal basal plane axes.

These magnetic wavevectors generate additional satellite reflections surrounding crystallographic Bragg peaks [31].

FIG. 1: (a) Real component of AC susceptibility phase diagram for Gd$_2$PdSi$_3$ showing the incommensurate (IC-I/II), previously proposed skyrmion (A) and paramagnetic (P) phases. (b) Crystal structure of Gd$_2$PdSi$_3$ with triangular planes of Gd ions (blue) between a honeycombed mix of Pd/Si (orange/cream). (c) Reciprocal space diagram showing the (300) structural peak and incommensurate satellites of interest, $\mathbf{\bar{q}}_n$. Orange arrow shows the relative direction of the outgoing x-ray wavevector ($\mathbf{\bar{k}}'$) to the magnetic wavevectors $\mathbf{\bar{q}}_n$. (d) Experimental set-up showing the magnetic field ($\mathbf{H}$)$||c$, the incoming $\pi$ polarization and the separation of the outgoing x-rays into $\pi$-$\pi$ or $\pi$-$\sigma$ components using an analyzer crystal.

Fig. 1(c) shows a set of three satellite reflections $\mathbf{\bar{q}}_{n,n=1,2,3}$, surrounding the (300) crystallographic Bragg peak, chosen for investigation as the magnetic modulation wavevector from $\mathbf{\bar{q}}_2$ is nearly collinear (to within $7^\circ$) with the scattered x-rays. This directional similarity together with the use of polarized x-rays allows the modulated moment within the basal plane to be broken down into components which are parallel and perpendicular to $\mathbf{\bar{q}}_n$. The experimental set-up utilizing polarized x-rays is shown in Fig. 1(d). In this horizontal scattering layout, both the incoming and outgoing x-ray wavevectors are within the crystallographic ab plane, and the magnetic field can be applied along the c direction. In our setup, intensity in the spin-flip channel ($\pi$-$\sigma$) arises from a magnetic modulation parallel to the incoming x-ray, which is in the ab plane, whereas the scattering in the non-spin-
flip $\pi-\pi$ channel arises from $c$-axis magnetic modulation and/or modulated charge order [31].

Figs. 2(a) and 2(b) show the zero-field energy scans of the scattered satellite intensity for each outgoing polarization at 8 K. In (a), the $\pi-\sigma$ channel (sensitive only to magnetic scattering) shows substantial intensity only close to the Gd L2 edge, and near-zero intensity everywhere, consistent with the expected large enhancement of the magnetic scattering cross-section on resonance [31]. By contrast, the $\pi-\pi$ channel in (b) (sensitive to both charge and magnetic scattering) shows significant intensity both at resonance and away from it. This intensity away from resonance cannot arise from magnetic scattering, as the non-resonant magnetic scattering component shown in (a) is over two orders of magnitude weaker than the peak at resonance, as expected for purely magnetic x-ray scattering [31], whereas in (b) the averaged off-resonance intensity is much larger at 20% of the resonant peak intensity. We note that it is impossible for this intensity away from resonance to arise from leak-through from the other polarisation channel, as we show this to be near zero in (a). This result shows that a substantial fraction of the incommensurate satellite $\pi-\pi$ intensity arises from charge scattering. As shown in Fig. 2(c) and 2(d), the small rocking curve width ($0.2^\circ$) of both the satellite and main Bragg peaks compared to their $3^\circ$ separation for both polarizations demonstrates that this scattering is well localized in reciprocal space. This demonstrates the presence of a low temperature CDW with the same wavevector as the magnetic structures.

We now turn our attention to the evolution of the magnetic and charge textures with an increasing magnetic field after zero field cooling. Fig. 3(a) shows the field-dependence of the AC susceptibility at 8 K, which displays clear peaks near 0.5 T and 1.05 T, highlighting magnetic transitions from the IC-I state to the A phase and the A phase to IC-II state, respectively. Fig. 3(b) shows the change in the magnitude of $\vec{q}_{\pi}$ obtained from the $\pi-\sigma$ (magnetic scattering only) channel and Fig. 3(c-e) shows the scattered intensity for both polarization channels for each of the three satellite peaks $\vec{q}_{n}$. The signature of the magnetic phase transition shown to occur at 0.5 T in (a) can also be seen by the variation of the magnetic wavevector in (b). That is, at low fields, the periodicity of the IC-I phase remains robust to deformations induced by the magnetic field and the wavevector remains constant. After the transition at 0.5 T, the wavevector now describes the magnetic texture within the A phase and is shown to continuously change in magnitude with further increases of magnetic field. This behaviour is in stark contrast to low-field magnetic helices in Ho and the incommensurate low-field helices within typical skyrmionic materials, where it has been shown the helical phase can continuously distort under the application of an applied field, while the spacing of the skyrmion lattice remains constant [32].

The intensities from each $\vec{q}_{n}$ in the $\pi-\sigma$ channels [orange curves in Fig. 3(c-e)] arise due to the interaction with an in-plane modulated magnetic moment. For each $\vec{q}_{n}$, we see an initial increase of intensity with increasing field, up to a maximum near the transition at 0.5 T, followed by a decrease to a non-zero value at the highest measured field of 0.95 T for $\vec{q}_{1,3}$, and near-zero for $\vec{q}_{2}$. We interpret this reduction to be caused by the moments uniformly aligning with the magnetic field, reducing the in-plane modulation at these wavevectors. In addition, we find that the $\pi-\sigma$ intensity from $\vec{q}_{3}$ peak is very small at all fields. This corresponds to a lack of a magnetic modulation along the direction of the magnetic wavevector [31], with the remnant intensity arising due to the small $7^\circ$ angle between the vectors. Our data is therefore consistent with in-plane magnetic moment oscillations for all fields lying perpendicular to the direction of the in-plane magnetic wavevector, in agreement with the findings of Kurumaji et al [4].

The intensities within the $\pi-\pi$ channels [blue curves on Fig. 3(c-e)] show further interesting features. At low fields, the intensities for the different satellites in the $\pi-\pi$ channel are not equal. Previously, residual strain arising from mounting the sample has been suggested to lift the degeneracy of magnetic and charge domains [33, 34]. This results in a higher population of one domain, explaining the increased scattering for $\vec{q}_{3}$. Upon increasing the field, the magnitude of $\vec{q}_{3}$ drops while the others remain constant, but ultimately results in nearly equal intensities for all three peaks at fields above 0.55 T.

Increasing the magnetic field past the phase transition at 0.5 T, we see a decrease of the intensity in the
FIG. 3: (a) Real component of the AC-susceptibility of Gd$_2$PdSi$_3$ as a function of magnetic field applied along c-axis. (b) Change in magnitude of the $\vec{q}_3$ magnetic wave vector for a similar field scan. (c–e) X-ray intensity separated into the $\pi-\pi$ (blue) and $\pi-\sigma$ (orange) channels as a function of applied field along the c-axis for each of the incommensurate satellites $\vec{q}_{in}$. The temperature for all measurements was 8 K. (f) Geometrically normalized $\pi-\sigma$ intensities for all $\vec{q}_{in}$ assuming the moment-modulation is perpendicular to $\vec{r}_3$. (g) Temperature dependence of the $\vec{q}_3$ satellite with polarisation separated intensity, $\pi-\pi$ (blue) and $\pi-\sigma$ (orange), normalised to the maximum scattered intensity.

$\pi-\pi$ channel, stabilizing at 0.75 T to a near-zero intensity. This field of 0.75 T does not match any of the magnetic phase transitions observed from AC susceptibility; in fact, it lies directly between the two transitions at the point where we would expect the magnetic A-phase to be most stable. It is therefore unlikely that this change in scattering results from a change in the magnetic state, reinforcing the picture that the bulk of the scattering in the $\pi-\pi$ channel is due to charge scattering, and not due to magnetism. Furthermore, a step-like feature in the magnetoresistance at 0.75 T has previously been reported [6], which matches the field at which we observe the $\pi-\pi$ scattering to fall off. Features in the magnetoresistance could arise from a change in the CDW state in the system, thus our observation of a collapse in the charge scattering at this field provides an explanation for this previously reported magnetoresistance feature.

Information about the relative magnetic domain populations is shown by the geometrically normalized $\pi-\sigma$ intensity, $I_n/(\vec{r}_n \cdot \vec{k}_{i,n})^2$, assuming the in-plane moment is perpendicular to the magnetic wavevector, in Fig. 3(f). We find equal scattering within errors for all fields in which the IC-I phase exists, which suggests that the previously mentioned non-equal $\pi-\pi$ scattered intensities at low field arise due to a different population of CDW domains. After the magnetic transition, the normalized intensities from the three wavevectors initially overlap at 0.55 T, and then continually drift apart under an increasing field up to 0.95 T. These features cannot arise from large-scale spin-reorientations, which would induce a large signal in the AC-susceptibility.

To further investigate the evolution of the magnetism and the CDW, we measured the temperature dependence of the scattered intensity from $\vec{q}_3$ at zero field and present it in Fig. 3(g), which has been normalized such that the low-temperature intensity is the same for both polarization channels. The different polarisation channels show that the onset of the CDW (blue/$\pi-\pi$) occurs at an appreciably lower temperature than the incommensurate magnetism (red/$\pi-\sigma$), suppressed by ~2 K. A similar feature has been seen in the related material Gd$_3$Ru$_4$Al$_{12}$, and was attributed to the evolution of a magnetic helix (which has a c-axis modulation) from a 1-dimensional transverse spin density wave (which has no c-axis component) in Ref. [35]. However, their analysis was ambiguous due to the possible presence of a CDW being indistinguishable at resonance to the supposed magnetic c axis component.

Instead, we suggest that at least in Gd$_2$PdSi$_3$, the CDW appearing at lower temperatures than the magnetic textures is consistent with the CDW being induced by the coupling between the itinerant electrons and the incommensurate local moments [28], rather than a direct interaction with the Fermi-surface nesting. The spin-coupling to a CDW and the large asymmetry between the periodic out-of-plane and in-plane magnetic modulations are in contrast to the previously suggested magnetic textures in this material. Recent local-force calculations have suggested that the zero-field incommensurate structures in Gd$_2$PdSi$_3$ are stabilized by inter-orbital frustration due to ferromagnetic exchange coupling between the itinerant Gd-5$d$ electrons and the antiferromagnetic coupling between local Gd-4$f$ magnetic moments [25]. The study also finds that the out-of-plane spin correlation is weak, further suggesting the magnetic textures only develop incommensurate coplanar modulation within the basal ab plane, and were unable to stabilize the purported non-coplanar magnetic skyrmions using this frustration mechanism [4]. Our measurements are in strong agreement with these theoretical findings.

We therefore suggest alternative, coplanar magnetic textures for the IC-I and A phases in Gd$_2$PdSi$_3$. These are shown in Fig. 4(a-c). In (a), we show a single-q transverse spin density wave (T-SDW) with the magnetic moment modulation within the ab plane. This state is degenerate with three symmetry-equivalent incommensurate wavevectors along the a, b-type crystal directions, and can form multiple macroscopic single-q domains. In
FIG. 4: (a) One-dimensional axis-modulated magnetic structure. The magnetic moment is periodically oscillating in a direction perpendicular to its wavevector and in-plane. (b, c) XY-vortex lattice composed from three superimposed one-dimensional structures oriented $120^\circ$ apart with relative intensities taken from Fig 3(f) at 0.55 T and 0.95 T respectively.

(b), we show a hexagonal lattice of coplanar amplitude modulated XY-vortices (XY-V), which can be deformed into a distorted vortex state (DXY-V), an example of which is shown in Fig. 4(c). The XY-V and DXY-V spin textures were generated using 3 superimposed T-SDW with wavevectors oriented $60^\circ$ apart, with equal weighting of each wavevector for XY-V. For the DXY-V in (c), the weighting of each wavevector was given by the geometrically normalised intensities at 0.95 T shown in Fig. 3(f).

The lack of phase information available in diffraction techniques typically makes determining the exact form of the magnetic textures non-trivial. Specifically, the signatures of the XY-V phase and 3 equally populated domains of T-SDW are indistinguishable with the set of $\vec{q}$, we investigated. Therefore we are unable to categorically assign these coplanar spin textures to either the IC-I or A-phase, and instead propose two possible scenarios.

**Scenario one:** The IC-I phase is comprised of 3 equally populated domains of T-SDW. Applying a magnetic field in the $c$-direction does not break the degeneracy of these domains, and we therefore expect near equal domain populations up to 0.5 T as shown in Fig. 3(f). Each T-SDW is coupled to a CDW with the same wavevector. This coupling would also explain the remarkable observation that the low-field magnetic wavevector does not deform with increasing field, as the coupling with CDW may stabilize the IC-I spin texture via cooperative organization [29], preventing distortions that would change the in-plane incommensurate ordering wavevector. The A-phase is initially comprised of a hexagonal lattice of XY-V, which deforms with an increasing field to an DXY-V. The lack of a c-axis magnetic modulation in coplanar textures prevents any solid angle spanned by neighbouring moments, inducing zero Berry phase required for the topological Hall effect [36]. However, it has been shown that gradients of the magnetization contribute to a theoretically greater signal known as the chiral Hall effect, which can even dominate the topological Hall effect that is associated with topologically-nontrivial noncoplanar spin textures such as skyrmions. [37]. This phenomena would be present in an XY-V lattices due to spontaneous symmetry breaking of vortex helicity inducing local chirality as well large gradients in magnetization, suggesting the origin of the large Hall-effect lies with the magnetic texture.

**Scenario two:** The IC-I phase is a hexagonal lattice of XY-V, which are not coupled to CDWs, as evidenced by the contrasting field-driven behaviour of the CDW and IC-I diffracted intensities in Fig. 3(c-f). The topological properties of the XY-V prevents lattice distortions which would result in a change of the magnitude of the wave vector, which would require a change the number of vortices. At fields above 0.5 T, 3 domains of T-SDW arise. Different domain populations may arise due to a field induced strain arising from the torque of the large Gd$^{3+}$ moment. This would suggest that the large Hall signal is due to the non-trivial electronic behaviour and the field-driven evolution the Fermi-surface, as Fermi-surface gaps have been shown to exhibit similar Hall signals [38]. A further study of the interplay between the electronic and magnetic states within the system, together with local-probe measurements, would have the potential to fully elucidate the magnetic textures and the origin of the large Hall signal.

In conclusion, we investigated the x-ray scattering from incommensurate structures in Gd$_2$PdSi$_3$ at various temperatures and fields. Energy scans highlight the presence of a CDW with the same incommensurate wavevector as the low-field IC-I magnetic state. Using polarisation-dependent scattering, we show that the onset of the CDW scattering occurs at lower temperatures than IC-I phase magnetic scattering, which suggests that the coupling between the local-moment order and the itinerant electrons drives the formation of the CDW. Intensity from the CDW is reduced to zero at 0.75 T, a field at which the AC susceptibility is at a mimima. Thus the polari-
sation channel previously used to infer a magnetic modulation in the $c$-direction is in fact dominated by charge scattering effects. We find that the IC-I and A-phases are highly-anisotropic, in agreement with recent theoretical findings [25]. Instead we suggest these phases are comprised of highly anisotropic, coplanar spin vortices and transverse spin-density waves, instead of the previously claimed non-coplanar skyrmions [4]. Our results call for urgent reinterpretation of recent results on this compound as a candidate frustrated skyrmion host.

The authors thank Diamond Light Source for time on Beamline I16 under Proposal MT20056-I. This work was financially supported by the UK Skyrmion Project EP/S021284/1. M. N. Wilson acknowledges the support of the Natural Sciences and Engineering Research Council of Canada (NSERC). M. Schmidt acknowledges the support of the Slovenian Research Agency under Project No. Z1-1852.

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