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New magnetic phase of the chiral skyrmion material Cu$_2$OSeO$_3$

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The lack of inversion symmetry in the crystal lattice of magnetic materials gives rise to complex noncollinear spin orders through interactions of a relativistic nature, resulting in interesting physical phenomena, such as emergent electromagnetism. Studies of cubic chiral magnets revealed a universal magnetic phase diagram composed of helical, conical, and skyrmion crystal phases. We report a remarkable deviation from this universal behavior. By combining neutron diffraction with magnetization measurements, we observe a new multidomain state in Cu$_2$OSeO$_3$. Just below the upper critical field at which the conical spiral state disappears, the spiral wave vector rotates away from the magnetic field direction. This transition gives rise to large magnetic fluctuations. We clarify the physical origin of the new state and discuss its multiferroic properties.

INTRODUCTION

Chiral magnets show a variety of periodically modulated spin states—spirals (1, 2), triangular and square arrays of skyrmion tubes (3–8), and a cubic lattice of monopoles and antimonopoles (9)—which can be viewed as magnetic crystals of different symmetries and dimensionalities. These competing magnetic superstructures show high sensitivity to external perturbations, allowing the control of phase boundaries with applied electric fields and stresses (10, 11). The nontrivial topography of multiply periodic magnetic states gives rise to emergent electromagnetic fields and unconventional spin, charge, and heat transport (12–16). The stability and small size of magnetic skyrmions as well as low spin currents required to set them into motion paved the way to new prototype memory devices (17–21).

Recent studies of chiral cubic materials hosting skyrmions, such as the itinerant magnets, MnSi and FeGe, and the Mott insulator, Cu$_2$OSeO$_3$, showed that they exhibit the same set of magnetic states with one or more long-period spin modulations and undergo similar transitions under an applied magnetic field (22). This universality is a result of noncentrosymmetric cubic lattice symmetry and the hierarchy of energy scales (2, 23, 24). The transition temperature $T_c$ is determined by the ferromagnetic (FM) exchange interaction $J$. The relatively weak antiferromagnetic Dzyaloshinskii-Moriya (DM) interaction, with the strength $D$ proportional to the spin-orbit coupling constant $\lambda$, renders the uniform FM state unstable toward a helical spiral modulation (1, 25).

It determines the magnitude of the modulation wave vector $Q$ and the value of the critical field $H_{c2}$ above which the spiral modulation is suppressed. In contrast to low-symmetry systems (26, 27), the DM interaction in cubic chiral magnets does not impose constraints on the direction of the spiral wave vector (2). The direction of the wave vector is controlled by the applied magnetic field and magnetic anisotropies of higher order in $\lambda$. In the helical spiral phase observed at low magnetic fields, magnetic anisotropies pin the direction of $Q$ either along one of the cubic body diagonals, as in MnSi, or along the cubic axes, as in FeGe or Cu$_2$OSeO$_3$. The competition between the Zeeman and magnetic anisotropy energies sets the critical field $H_{c1}$ of the transition between the helical and conical spiral states, above which $Q$ is parallel to the applied magnetic field. In the multiply periodic skyrmion crystal state, the spiral wave vectors are perpendicular to the field direction, which is favored by the nonlinear interaction between the three helical spirals.

Here, we report a remarkable deviation from this well-established universal behavior. By small-angle neutron scattering (SANS) and magnetic measurements, we observe a new low-temperature magnetic phase of Cu$_2$OSeO$_3$. At relatively high magnetic fields, $Q$ tilts away from the magnetic field vector, $H$, when this is directed along the [001] crystallographic direction favored by anisotropy at zero field. This transition occurs where it is least expected—at $H$ close to $H_{c2}$, where the dominant Zeeman interaction favors $Q \parallel H$ and at low temperatures where thermal spin fluctuations that can affect the orientation of $Q$ are suppressed. The reorientation of the spiral wave vector is accompanied by strong diffuse scattering, reminiscent of the pressure-induced partially ordered magnetic state in MnSi (28). The instability of the conical spiral state at high applied magnetic fields can be considered as a re-entrance into the helical state, although $Q$ in the “tilted conical spiral” state is not close to high-symmetry points. The new phase of Cu$_2$OSeO$_3$ is sensitive to the direction of the applied magnetic field: For $H \parallel (110)$, no tilted spiral state is observed. Instead, we find that the helical-to-conical spiral transition splits into two transitions occurring at slightly different magnetic fields.

We show theoretically that the tilted spiral state originates from the interplay of competing anisotropic spin interactions, which is generic to chiral magnets and may be important for understanding the structure of metastable skyrmion crystal states (7, 29). This interplay is particularly strong in Cu$_2$OSeO$_3$ owing to the composite nature of spin of the magnetic building blocks (30). The transition to the new state in multiferroic Cu$_2$OSeO$_3$ should have a strong effect on the magnetically induced electric polarization. It should also affect the spin Hall magnetoresistance (31) and modify the spin wave spectrum.
RESULTS

First hints for the existence of the new phase came from the anomalous field dependence of the magnetization, $M$, and the ac magnetic susceptibilities, $\chi'$ and $\chi''$, shown in section S1. A direct confirmation was provided by SANS, which probes correlations perpendicular to the incoming neutron beam wave vector $k_i$. For this reason, we performed our measurements in two crystallographic orientations and for each orientation in two complementary configurations, $H \perp k_i$ and $H \parallel k_i$, thus providing a full picture of the effect of the magnetic field on the magnetic correlations.

A selection of patterns obtained at $T = 2$ K is shown in Fig. 1 (A and B) for $H \parallel [110]$ and in Fig. 1 (C and D) for $H \parallel [001]$. At zero field, the SANS patterns show four peaks along the diagonal directions in Fig. 1 (A and D) for $k_i \parallel [001]$ and two peaks along the horizontal axis in Fig. 1 (B and C) for $k_i \parallel [110]$. These are the magnetic Bragg peaks of the helical spiral state with wave vectors along the three equivalent (001) crystallographic directions.

At $\mu_0 H = 25$ mT, the scattered intensity vanishes for $H \parallel k_i$ (Fig. 1, B and D) because of the reorientation of the spiral wave vector along the magnetic field at the transition to the conical spiral phase. On the other hand, for $H \perp k_i$ and $H \parallel [110]$, Fig. 1A shows the coexistence of helical spiral and conical spiral peaks (additional weak peaks are attributed to multiple scattering). Thus, the helical-to-conical transition for $H \parallel [110]$ is not a simple one-step process. Reorientation first occurs in the helical spiral domain with the wave vector perpendicular to the field direction $Q \parallel [001]$. It is followed by a gradual reorientation of the wave vectors of the other two helical spiral domains. Upon a further increase of the magnetic field, the conical spiral peaks weaken in intensity and disappear at the transition to the field-polarized collinear spin state, which, for $H \parallel \langle 110 \rangle$, occurs above 75 mT.

The unexpected behavior, a signature of the new phase, is seen in the evolution of SANS patterns for $H \parallel [001]$ in Fig. 1 (C and D). For $H \perp k_i$ (Fig. 1C), the Bragg peaks broaden along the circles with radius $Q$ and eventually split into two well-defined peaks at 35 and 40 mT. This is surprising because, in this configuration, the two peaks along the horizontal axis correspond to the spiral with the wave vector parallel to both the magnetic field and the cubic axis. Thus, no reorientation is expected for the spiral domain favored by both the Zeeman interaction and magnetic anisotropy. In addition to the splitting of the Bragg peaks, in the complementary configuration of $H \parallel k_i$ shown in Fig. 1D, a broad ring of scattering develops well inside the circle with radius $Q$.

With increasing temperature, the splitting of the Bragg peaks becomes smaller and disappears at $\sim 50$ K, as shown in Fig. 2A. At $T = 10$ K, the scattered intensity on the circle with radius $Q$, when plotted against the azimuthal angle $\phi$, consists of two Gaussians peaks, labeled 1 and 2 (Fig. 2B). These are centered at two distinct angles, which vary with temperature, and their difference reaches 30° at $T = 2$ K (Fig. 2C). The integrated intensities depicted in Fig. 2D show that peak 2, which splits away from the conical spiral peak 1, is by far the more intense one. Its intensity goes through a maximum at $\sim 35$ K and then decreases at low temperatures, possibly because the optimum Bragg condition is not fulfilled any longer as the peak moves away from the magnetic field direction.

Our experimental findings are summarized in Fig. 3, which shows contour plots of the real and imaginary susceptibilities, $\chi'$ and $\chi''$, as well as the phase boundaries obtained by SANS. Close to $T_c$, the
Fig. 2. Temperature dependence of the tilt. (A) Temperature dependence of the SANS patterns and (B to D) the corresponding data analysis. The dashed circles on the SANS patterns are a guide to the eye and have a radius of $Q = 2\pi/l \sim 0.1$ nm$^{-1}$. The angular dependence of the scattered intensity along the circle with radius $Q$ is given in (B) for $T = 10$ K. The solid lines indicate a fit with two Gaussian peaks labeled 1 and 2 in the SANS patterns. The temperature dependence of the angular positions and integrated intensities of the two peaks is shown in (C) and (D), respectively.

Fig. 3. Phase diagrams from ac magnetic susceptibilities. Contour plots of $\chi'$ and $\chi''$ at a frequency of 10 Hz and phase boundaries obtained by SANS for $H \parallel [110]$ (A to C) and $H \parallel [001]$ (D to F). The units for $\chi'$ and $\chi''$ are $10^{-4}$ and $10^{-6}$ m$^3$/mol Cu, respectively. The helical, conical, A, tilted spiral (TS), and field-polarized (FP) phases are indicated in (C) and (F). The phase boundaries determined from the susceptibility are illustrated by the symbols and the dashed lines in (A), (B), (D), and (E). They correspond to the peaks of $\chi''$, with the exception of $\mu_0 H_{C2}$, which is defined by the inflection point of $\chi'$ versus $\mu_0 H$. These criteria are the same as in our previous study (39). At low fields, two lines $\mu_0 H_{C1}^{(1)}$ and $\mu_0 H_{C1}^{(2)}$ are identified below 50 K. Just below $\mu_0 H_{C2}$, a red dashed line denoted as $\mu_0 H_{CT}$ in (D) to (F) marks the onset of the “tilted spiral” state for $H \parallel [001]$. The phase boundaries determined from SANS are illustrated by the green symbols in (C) and (F). The shaded gray area just below $\mu_0 H_{C2}$ in (F) marks the region where the ring of scattering emerges for $H \parallel [001] \parallel k$. 
transition from the helical to the conical phase is marked by a single 
\(\mu_0H_{C1}(T)\) line, which, at low temperatures, evolves into two lines,
\(\mu_0H_{C1}^{(1)}(T)\) and \(\mu_0H_{C1}^{(2)}(T)\), derived from the two adjacent \(\chi''\) peaks
(see fig. S1, C and F, and the discussion in the Supplementary Materials).
The most prominent difference between the two field orientations
appears close to \(\mu_0H_{C2}\) below 30 K. In this field and temperature range,
clear maxima are seen in both \(\chi'\) and \(\chi''\) for \(H||[001]\). These define a new
line \(\mu_0H_{C2}(T)\) (red dashed line in Fig. 3, D to F), which shifts slightly to
lower magnetic fields with decreasing temperature.

The boundaries determined from the SANS measurements, shown in
Fig. 3 (C and F), are in excellent agreement with those derived from
susceptibility. Furthermore, it is remarkable that the shaded area in Fig.
3F, which marks the region where the ring of scattering shown in Fig.
1D emerges for \(H||[001]\) \(\parallel k_0\) coincides with the maxima of \(\chi'\) and \(\chi''\).

**DISCUSSION**

This remarkable behavior can be understood in terms of competing
magnetic anisotropies that are generic to cubic chiral magnets. This
competition is particularly tight in Cu2OSeO3, as explained below. Despite
the long history of studies of cubic chiral magnets (23, 24), the discussion of
anisotropic magnetic interactions in these materials remains,
to the best of our knowledge, incomplete.

The starting point of our approach is a continuum model with
the free energy per unit cell

\[
\varepsilon = \frac{1}{2} a^2 \sum_{i=x,y,z} \partial_i m \partial_i m + D a m \cdot \nabla \times m - a^3 \mu_0 MmH + \varepsilon_a
\]

(1)

where \(m\) is the unit vector in the direction of the magnetization, \(M\) is
the magnetization value, \(a\) is the lattice constant, and \(\varepsilon_a\) is
the magnetic anisotropy energy. There are five terms of fourth order
in the spin-orbit coupling, \(\lambda\), allowed by the P213 symmetry

\[
\varepsilon_a = C_1 a^2 \left[ \partial_x m_x \partial_x m_x + \partial_y m_y \partial_y m_y + \partial_z m_z \partial_z m_z \right]
+ C_2 a^2 \left[ \partial_z m_x \partial_y m_x + \partial_x m_y \partial_z m_y + \partial_y m_z \partial_x m_z \right]
- \left( \partial_z m_x \partial_y m_y + \partial_y m_z \partial_x m_z \right)
+ 2C_3 a^2 \left[ \partial_x m_y \partial_z m_y + \partial_y m_z \partial_x m_z \right]
+ \frac{1}{2} J a^4 \left[ \partial_x^2 m \cdot \partial_x^2 m + \partial_y^2 m \cdot \partial_y^2 m + \partial_z^2 m \cdot \partial_z^2 m \right]
+ K \left( m_x^4 + m_y^4 + m_z^4 \right)
\]

(2)

Their importance can be understood by substituting into Eq. 1 the
conical spiral Ansatz

\[
m = \cos \theta \epsilon_3 + \sin \theta \left[ \cos(Qx) \epsilon_1 + \sin(Qx) \epsilon_2 \right]
\]

(3)

where \(\theta\) is the conical angle and \(\{\epsilon_1, \epsilon_2, \epsilon_3\}\) are three mutually orthogonal unit vectors. If the magnetic anisotropy and Zeeman energies are
neglected, \(Qa = \frac{D}{2}\) is independent of the orientation of \(Q\). The applied
magnetic field favors \(Q||H\) with the conical angle given by 
\(\cos \theta = \frac{H_{C2}}{H}\)
where \(\mu_0MH_{C2}\) defines the critical field \(H_{C2}\).

The DM interaction originates from the antisymmetric anisotropic exchange
between Cu spins, which is the first-order correction to the
Heisenberg exchange in powers of \(\lambda; D \sim J_0^2\), where \(\zeta = \frac{3}{2}\), \(\Delta\) being the
typical electron excitation energy on Cu sites (25, 32). The first three
anisotropy terms in Eq. 2 result from the symmetric anisotropic exchange
between Cu ions and are proportional to the second power of
the spin-orbit coupling (25, 32); \(C_1 \sim J_0^2\). Since \(Qa = \frac{D}{2} = \zeta\), these
anisotropy terms are of the order of \(J_0^4\). The fourth term in Eq. 2 results
from the expansion of the Heisenberg exchange interaction in powers of
\(Qa\) and is also \(\sim J_0^4\). The last term in Eq. 2 has the form of the fourth-
order single-ion anisotropy allowed by cubic symmetry. In absence of
the single-ion anisotropy for Cu ions with \(S = \frac{1}{2}\), this term emerges at
the scale of the unit cell containing 16 Cu ions, which form a network
of tetrahedra with \(S = 1\). This last term appears either as a second-
order correction to the magnetic energy in powers of the symmetric
anisotropic exchange or as a fourth-order correction in the antisymmetric exchange, both \(\propto \zeta^4\). The intermediate states are excited states of Cu
tetrahedra (34) with energy \(\sim J^2\) rather than the electronic excitations of
Cu ions with energy \(\Delta\). As a result, the last term is also \(\sim J_0^4\). Therefore,
the magnetic block structure of Cu2OSeO3 makes all anisotropy terms in
Eq. 2 comparable, which can frustrate the direction of \(Q\). We note
that the ferrimagnetic nature of Cu2OSeO3 does not play an important
role: For the magnetic fields considered here and the large spin gap in
the spin tetrahedron of about 275 K (30, 33), the tilt angle between the
two magnetic sublattices should be less than \(10^{-4}\) and can be neglected.

Another important point is that the direction of \(Q\), favored by
a magnetic anisotropy term, may vary with the strength of the applied
magnetic field because it depends on the conical angle, \(\theta = \theta(H)\). Figure
S6 shows the \(\theta\) dependence of the fourth-order effective anisotropy,
\(K_{eff} = KB(\theta)\), which is negative for small \(\theta\) and for \(\theta \approx \frac{\pi}{2}\), stabilizing
the helical spiral with \(Q||[001]\), as it is the case for Cu2OSeO3 at zero field.
However, for intermediate values of \(\theta\), \(K_{eff}\) is positive and the preferred
direction of \(Q\) becomes \(\langle 111 \rangle\). In this interval of \(\theta\), spins in the conical
spiral with \(Q||[001]\) are closer to the body diagonals than to cubic axes,
which makes this wave vector direction unfavorable (see section S4 for
more details).

This effect gives rise to local minima of the conical spiral energy in
\(Q\)-space. If only the fourth-order anisotropy is taken into account, the
global energy minimum for \(H||[001]\) is still at \(Q||H\). However, addi-
tional anisotropy terms can turn these local minima into global ones,
such that, in some magnetic field intervals, the tilted conical spiral
becomes the ground state. Figure 4 (A to F) shows the false color plot of
the conical spiral energy as a function of \(Q\) for several values of the
dimensionless magnetic field, \(h = H/H_{C2}\), applied along the \langle 001 \rangle
direction. For simplicity, only two anisotropy terms are nonzero in this
calculation: \(K = \frac{K_1}{a^2} = -0.19\) and \(\gamma = \frac{\gamma_1}{a^2} = -0.1\). In zero field (Fig. 4A),
there are three energy minima along \langle 001 \rangle, that is, along the \langle 001 \rangle,
\langle 010 \rangle, and \langle 001 \rangle directions, corresponding to three degenerate
conical spiral domains in Cu2OSeO3. For \(h = 0.2\) (Fig. 4B), the helical spiral
states with \(Q\) along the \langle 001 \rangle and \langle 010 \rangle directions are metastable.
For \(h = 0.3\) (Fig. 4C), only the minimum with \(Q||[001]\) exists, corresponding
to the conical spiral state. For \(h = 0.6\) (Fig. 4D), the conical spiral with \(Q||[001]\)
is unstable and there appear four new minima, corresponding to
four domains with \(Q\) tilted away from the magnetic field vector along
the \([\pm 1, \pm 1, 0]\) directions, as can be seen more clearly in Fig. 4F showing
the energy sphere seen from above. The relative energy changes in this

part are \( \sim 10^{-2} \), implying large fluctuations of the spiral wave vector, which can explain the diffuse scattering shown in Fig. 1D. Finally, Fig. 4E shows the field-polarized state at \( h = 1 \).

Figure 4 (G to L) shows the \( Q \) dependence of the conical spiral energy in several magnetic fields along the [110] direction, calculated for the same values of parameters as those in Fig. 4 (A to F). As the magnetic field increases, the helical spiral states with \( Q \) along the [001] and [010] directions merge into a single state with the wave vector parallel to \( H \), and the state with \( Q \parallel [001] \) ultimately disappears (see Fig. 4, G to I). This gives rise to the two-step transition from the helical to the conical phase observed experimentally. For this field direction, the multidomain tilted spiral state does not appear and there is only one global minimum with \( Q \parallel [110] \) for \( H > H_{C2} \). Nevertheless, one can see the strong vertical elongation of the energy contours in Fig. 4 (H to J), which is a result of the competition with the \( Q \parallel [11 \pm 1] \) states. At \( h = 1 \), the elongation changes from vertical to horizontal (see Fig. 4, K and L).

These conclusions drawn using the variational approach are confirmed by exact energy minimization of Eq. 1 including two competing anisotropy terms: the fourth-order anisotropy with \( \kappa = -0.2 \) and the anisotropic exchange \( \gamma_1 \). The case of \( H \parallel [110] \) is treated in section S5, which explains the two lines \( \mu_0 H_{C1}^{[110]} \) and \( \mu_0 H_{C2}^{[110]} \), for the transition from the helical to the conical spiral state, shown in Fig. 3 (A to C). For \( H \parallel [001] \), the field dependence of the angle, \( \phi \), between \( Q \) and the [001] cubic axis, which describes the tilt of \( Q \) toward the [111] directions, is shown in Fig. 5A. The tilted spiral state appears when \( |\gamma_1| \) exceeds a critical value, which is slightly lower than 0.1 for \( \kappa = -0.2 \). When the magnetic field increases, the tilt angle reaches its maximal value, \( \phi_{\text{max}} \), marked by the empty circles in Fig. 5A, and then decreases to 0. As shown in Fig. 5B, \( \phi_{\text{max}} = 0 \) for \( |\gamma_1| \leq 0.1 \). Thus, the state with \( Q \parallel [110] \) is stable at low anisotropies. However, as the exchange anisotropy increases, an intermediate state occurs, and finally, for \( |\gamma_1| \geq 0.28 \), the state with \( Q \parallel [111] \) is stabilized even at zero magnetic field.

In our diffraction experiment, we do not observe all four tilted spiral domains, which is likely related to a small misalignment of the sample: Because of a weak dependence of the spiral energy on \( Q \), even a tiny deviation of \( H \) from the [001] direction leads to the selection of one of the four domains. This suggests that the domain structure of the tilted state can also be controlled by an applied electric field using the multiferroic nature of \( \text{Cu}_2\text{OSeO}_3 \) (35–37). The electric polarization induced by the tilted spiral with the spin rotation axis \( I = (\frac{1}{\sqrt{2}} \sin \alpha, \frac{1}{\sqrt{2}} \sin \alpha, \cos \alpha) \) is given by

\[
\langle P \rangle = \lambda \frac{(3\cos^2 \theta - 1)}{4\sqrt{2}} \left( \sin 2\alpha, \sin 2\alpha, \frac{1}{\sqrt{2}}(1 - \cos 2\alpha) \right)
\]

For small anisotropies, \( \alpha \) is close to the tilt angle of \( Q \), \( \phi \) (a more precise relation between \( \alpha \) and \( \phi \) and the derivation of Eq. 4 can be...
found in section S6). Since \( \phi \) does not exceed 30\(^{\circ} \), the induced electric polarization is almost normal to the applied magnetic field \( H \parallel [001] \). The conical spiral with \( \alpha = 0 \) induces no electric polarization.

To summarize, we found a new high-field multidomain magnetic state that intervenes between the conical spiral and field-polarized phases and is stable in a broad temperature range. This major deviation from universal behavior has important consequences for the field of chiral magnetism, for example, the partial magnetic order observed in MnSi under pressure \( \langle 111 \rangle \), spin Hall magnetoresistance, and anomalies in the spin wave spectrum. The spiral tilt can also give rise to new topological magnetic defects, such as skyrmions, with interesting static and dynamic properties. Our theoretical model that takes into account all competing magnetic anisotropies is in a semiquantitative agreement with experiment. For particular directions of the magnetic field, this competition may stabilize skyrmions \( (38) \) or new tilted phases. Although the spiral wave vector tilts toward the \( (111) \) directions, the energy difference with \( Q \) along the \( (100) \) directions is relatively small. Therefore, one may envisage the formation of a superstructure with smooth rotation from one spiral domain to another, that is, a conical spiral composed of tilted spirals.

**MATERIALS AND METHODS**

Magnetization and magnetic susceptibility measurements were performed on two single crystals of \( \text{Cu}_2\text{OSeO}_3 \) with dimensions of \( \sim 1 \text{ mm} \times 1 \text{ mm} \times 1 \text{ mm} \) grown at the Zernike Institute for Advanced Materials. One crystal was oriented with a \( (001) \) axis vertical, while the other one was oriented with a \( (110) \) axis vertical. A third single crystal with dimensions of \( \sim 3 \text{ mm} \times 3 \text{ mm} \times 4 \text{ mm} \) grown at the Max Planck Institute for Chemical Physics of Solids was used for the neutron scattering measurements. This sample was oriented with the \( [110] \) crystallographic axis vertical. All crystals were prepared by chemical vapor transport method, and their quality and structure were checked by x-ray diffraction.

Magnetization and magnetic susceptibility were measured with an MPMS 5XL SQUID using the extraction method. For the determination of the magnetization, a static magnetic field, \( \mu_0 H \), was applied along the vertical direction. The real and imaginary parts of the magnetic ac susceptibility, \( \chi' \) and \( \chi'' \), were measured by adding to \( \mu_0 H \) a vertical drive ac field, \( \mu_0 H_{ac} \), with an amplitude of 0.4 mT. The frequency of \( H_{ac} \) was varied between 0.1 and 1000 Hz.

The SANS measurements were performed on the instruments PA20 of the Laboratoire Léon Brillouin and GP-SANS of the Oak Ridge National Laboratory using neutron wavelengths of 0.6 and 1 mm, respectively. At both instruments, the magnetic field was applied either parallel or perpendicular to the incoming neutron beam wave vector \( k_i \) using a horizontal magnetic field cryomagnet. The orientation of the crystal axes with respect to \( k_i \) and to the magnetic field was varied by rotating the sample in the cryomagnet. The SANS patterns were collected for \( H \parallel [110] \) and \( H \parallel [001] \) and, in each case, for \( H \perp k_i \) and \( H \parallel k_i \), by rotating both the sample and the magnetic field through 90\(^{\circ}\) with respect to \( k_i \). Measurements at 70 K, where the magnetic scattering is negligible, were used for the background correction of the SANS patterns.

All measurements were performed after zero field cooling the sample through the magnetic transition temperature, \( T_c \), down to the temperature of interest. The magnetic field was then increased stepwise. The applied magnetic field, \( \mu_0 H_{ext} \), was corrected for the demagnetizing effect to obtain the internal magnetic field, \( \mu_0 H_{int} \) (in SI units)

\[
H_{int} = H_{ext} - NM
\]

where \( N = \frac{1}{3} \) is the demagnetization factor for our (nearly) cubic shape samples. The demagnetizing field correction also modifies the values of the magnetic susceptibility

\[
\chi_{int} = \frac{\chi_{ext}}{1 - \mu_0 N \gamma' \chi_{ext}}
\]

**SUPPLEMENTARY MATERIALS**

Supplementary material for this article is available at http://advances.sciencemag.org/cgi/content/full/4/9/eaat7323/DC1

Section S1. Temperature and magnetic field dependence of the magnetization and susceptibility

Section S2. Frequency dependence of the ac susceptibility

Section S3. Phase boundaries determined from SANS

Section S4. Effect of magnetic anisotropy on the direction of the spiral wave vector

Section S5. Numerical studies of spiral reorientation processes in the presence of competing anisotropies

Section S6. Electric polarization

Fig. S1. Magnetic properties of \( \text{Cu}_2\text{OSeO}_3 \).

Fig. S2. Temperature dependence of the magnetization.

Fig. S3. \( \chi' \) and \( \chi'' \) as a function of magnetic field.

Fig. S4. Frequency dependence of \( \chi' \) and \( \chi'' \) at \( T = 5 \text{ K} \).

Fig. S5. Magnetic field dependence of the SANS intensity at 6 K.

Fig. S6. Effective anisotropy.

Fig. S7. Spiral reorientation for \( H \parallel [001] \).

Fig. S8. Spiral reorientation for \( H \parallel [110] \).

Fig. 50. Field-induced reorientation of the spiral wave vectors for \( H \parallel [001] \) and for the fourth-order anisotropy \( \kappa = -0.2 \) and the anisotropic exchange \( \gamma > 0 \).

Reference (40)

**REFERENCES AND NOTES**


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New magnetic phase of the chiral skyrmion material Cu$_2$OSeO$_3$

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