Meron crystals in chiral magnets

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Emergent stable topological structures are ubiquitous in condensed matter, high energy physics and materials science, as well as in atomic, molecular and optical physics. Prominent examples include domain walls, vortices, dislocations, disclinations and skyrmions. Skyrmions are vector field (e.g. spin) textures, which wrap a sphere once, and can be viewed as emergent mesoscale particles. These textures have been observed in quantum Hall states of two-dimensional electron gases, chiral nematic liquid crystals, and Bose-Einstein condensates, among others. The recent discovery of skyrmions in chiral magnets, e.g. MnSi, has triggered enormous interest because of their huge potential for spintronics applications. Unlike magnetic domain walls, skyrmions can be manipulated with very small electric currents, thus rendering them as prime candidates for novel information storage devices with much lower power consumption. Here we study ultrathin chiral magnets with easy-plane anisotropy and find that the skyrmion lattice evolves continuously into a meron-antimeron crystal. Because macroscopic functionality depends strongly on the topology of the emergent mesoscale particles, the stabilization of meron crystals may lead to novel macroscopic responses relevant to magnetic sensing and storage.

Skyrmion crystals [1][2] have been recently observed in chiral magnets. [3][4] Each skyrmion consists of a spin configuration that wraps the sphere once over a length scale that ranges from 10 to 100 nm [see Fig. 1(a)]. This property is quantified by the topological skyrmion charge $Q = \frac{1}{4\pi} \int d^2r \cdot (\partial_t \mathbf{n} \times \partial_x \mathbf{n}) = \pm 1$ [$\mathbf{n}(\mathbf{r})$ is a unit vector describing the direction of the spin with spatial coordinates $\mathbf{r}$]. The skyrmion texture affects the orbital motion of conduction electrons (coupled to the moments via Hund’s exchange), which experience an emergent magnetic flux proportional to the skyrmion charge: $\Phi = Q \Phi_0$ ($\Phi_0 = \hbar c/e$ is the flux quantum). The emergent magnetic field experienced by the conduction electrons is enormous: 100 T for a skyrmion size of 10 nm. This phenomenon leads to the so-called topological Hall effect. [8][9] However, most of the excitement generated by the recent discovery of mesoscale skyrmions in chiral magnets arises from their transport properties. Skyrmions can be driven by spin polarized currents in metals [10][12] or by magnon currents in insulators. [13][15] A very attractive feature is that the currents required to move skyrmions in chiral magnets are $10^5$ to $10^6$ times smaller than the depinning currents of magnetic domain walls. [10][12] making skyrmions prime candidates for novel information storage devices with substantially reduced power consumption. [16]

By splitting a skyrmion into two halves, we obtain a meron and antimeron (see Fig. 1). We will show here that this effect stabilizes meron crystals for strong enough easy-plane anisotropy and to explain the broader implications of this result.

Chiral magnets exhibit helical magnetic ordering at low enough temperatures. The skyrmion crystals emerge only in a small region of the magnetic field-temperature phase diagram of bulk chiral magnets. [3] In contrast, this exotic phase is stable over a wide range of the temperature-magnetic field phase diagram for thin films. [4][5] The key point is that spin anisotropy becomes important when the sample thickness is reduced below a certain length. [20][21] For instance, the easy-axis anisotropy of FeGe increases for thinner films, [21] which stabilizes the skyrmion phase in a broader range of temperature-magnetic region [22][23]. While spin anisotropy is relevant to stabilize the skyrmion crystal relative to the simple conical state, it is clear that a strong anisotropy is incompatible with the non-coplanar nature of the skyrmion texture. The inclusion of easy-plane anisotropy should push the spins away from the north and south poles towards the equator of the skyrmion (see Fig.1). We will show here that this effect stabilizes meron crystals for strong enough easy-plane anisotropy and that a sequence of different non-coplanar magnetic orderings appears in between the skyrmion and the meron crystal phases.

The Hamiltonian

$$\mathcal{H} = d \int d^2r \left[ \frac{J_{ex}}{2a} (\nabla \mathbf{n})^2 + \frac{D}{a^2} \mathbf{n} \cdot (\nabla \times \mathbf{n}) - \frac{B_z}{a^2} n_z + \frac{A}{2a^3} n_z^3 \right],$$

(1)
describes a 2D thin film of a chiral magnet with easy-plane anisotropy. [11][24] Here $d$ is the thickness, $a$ is the lattice constant, $J_{ex}$ is the ferromagnetic exchange and $D$ is the Dzyaloshinskii-Moriya (DM) interaction generally present in magnets without inversion symmetry. $B_z$ is the external magnetic field normal to the film and $A > 0$ is the easy-plane anisotropy. The film is assumed to be uniform along the $z$ direction. In the absence of magnetic field and anisotropy ($B_z = 0$ and $A = 0$), the DM interaction favors non-collinear orderings and therefore twists the uniform spin configuration
is

Lower panel: A skyrmion covers the sphere, whereas a meron covers only half the sphere, meaning that the topological charge of a skyrmion is $|Q| = 1$ and that of a meron is $|Q| = 1/2$. The regions where the spins are parallel and antiparallel to the applied field are indicated with red and blue colors respectively. Spins are parallel to the plane in the green region.

Figure 1. (color online) Upper panel: Dissociation of a skyrmion into a meron and antimeron pair in the presence of easy plane anisotropy. Lower panel: A skyrmion covers the sphere, whereas a meron covers only half the sphere, meaning that the topological charge of a skyrmion is $|Q| = 1$ and that of a meron is $|Q| = 1/2$. The regions where the spins are parallel and antiparallel to the applied field are indicated with red and blue colors respectively. Spins are parallel to the plane in the green region.

into a spiral with wave length $\lambda = J_{\text{ex}} a / D$. For $A = 0$, the triangular skyrmion crystal is stabilized in an intermediate magnetic field region, $0.2D^2 / J_{\text{ex}} < B_c < 0.8D^2 / J_{\text{ex}}$, and the spins become fully polarized above the saturation field. [24] For typical chiral magnet thin films, the skyrmion phase is stable in the magnetic field region $40 \text{ mT} < B_c < 80 \text{ mT}$ at low temperatures. [4] To find the different stable phases as a function of increasing easy-plane anisotropy, $A$, we simulate the dynamics of the spin field $\mathbf{n}(\mathbf{r})$ by numerically solving the Landau-Lifshitz-Gilbert equation with a noise term. The $T = 0$ state is obtained by slowly annealing the system from the paramagnetic state. [See Methods Section]

The evolution of the low temperature spin texture as a function of $A$ is depicted in the left column of Figs. 2 and 3. The middle column of Fig. 2 and the right column of Fig. 3 show the corresponding skyrmion density $q(\mathbf{r}) = \mathbf{n} \cdot (\partial_x \mathbf{n} \times \partial_y \mathbf{n}) / (4\pi)$. For small $A$, we recover the expected triangular skyrmion crystal. The spins are anti-aligned with the magnetic field at the center of each skyrmion, while they remain parallel to the field in the background region between skyrmions. As the easy-plane anisotropy increases, the skyrmion size increases because of a natural expansion of the equatorial region favoured by $A$. As a consequence of this expansion, skyrmions start overlapping with each other, i.e., the spins no longer wrap the full sphere ($Q < 1$) because the skyrmion size becomes larger than the separation between neighbouring skyrmions. The growth of the equatorial regions of three skyrmions that close a triangular loop [Fig. 2(d)], or a square loop [Fig. 2(g)], generates a vortex with an opposite winding number (antivortex) at the center of the loop. This antivortex is actually the precursor of an antimeron because the spins in the core region remain canted along the field direction but $|Q| < 1/2$ ($Q > 0$). We also note that the continuous suppression of the charge $Q$ of the blue “particles” of Figs. 2 and 5 implies that the former skyrmions ($Q = -1$) of Fig. 2(a) evolve continuously into the antimerons ($Q = -1/2$) of Figs. 3(a) and 3(c).

As it is shown in the different panels of Fig. 2, the precursors of the merons (blue particles) and antimerons (red particles) form two different crystals. For $A = 1.65D^2 / J_{\text{ex}}$, the crystal consists of a triangular lattice with red particles in one sublattice and blue particles in other two [see Figs. 2(d) and 2(e)]. The second crystal becomes stable for $2.55D^2 / J_{\text{ex}} < A < 2.85D^2 / J_{\text{ex}}$ and consists of two interpenetrating square lattices of blue and red particles [see Figs. 2(g) and 2(h)]. Interestingly, if we think of the blue and red particles as disks of different radii, these two crystals are the only compact packings [25] that satisfy the following rule: two small disks cannot be tangent to each other. This rule arises from an energetic consideration: spins near the point of contact of two disks of the same type (chirality) have to be parallel to the hard axis. The triangular and square crystals of Fig. 2(d) and 2(g) become compact packings when the ratio between the radii of the red and blue particles is equal to $r_r / r_b = 2\sqrt{3}/3 - 1$ and $r_r / r_b = \sqrt{2}-1$, respectively. [25] Because $r_r / r_b$ increases with $A$, it is natural to expect a transition from the triangular [Fig. 2(d)] to the square [Fig. 2(g)] interpenetrating crystals.
Figure 2. (color online) Evolution of the spin texture as a function of easy-plane anisotropy. Left column is the spin texture with color denoting the spin component along the $z$ direction while the arrows represent the in-plane spin component. Middle column is the corresponding topological charge density $q(r)$. From top to bottom, $A = 0.15D^2/J_{ex}$, $A = 1.65D^2/J_{ex}$, and $A = 2.55D^2/J_{ex}$. Note the transition from a triangular lattice to a square lattice with increasing anisotropy. The upper insets in (a, d, g) are the Fourier transform of $n_z$ after subtracting the uniform component, which can be measured by neutron scattering, while the lower insets are enlarged views. (c) and (f) show the corresponding wrapping of spin textures on the surface of a sphere. (i) illustrates why the chirality of the meron at the center is opposite to the surrounding merons. Here $B_z = 0.7D^2/J_{ex}$ and $L_x \times L_y = 60 \times 60 (J_{ex}a/D)^2$. The results are qualitatively the same for other magnetic fields.

As the easy-plane anisotropy is further increased, the particles separate from each other because of the continuous growth of equatorial region: now the background consists of a coplanar spin ordering in the $xy$-plane (see Fig. 3) and the former skyrmions or blue particles have evolved into merons with topological charge $Q = -1/2$. At the center of the $Q = -1/2$ meron loops (blue particles), there is always a $Q = 1/2$ antimeron (see Fig. 3). Spins have to come out of the plane in
Figure 3. (color online) Same as Fig. 2 but for higher values of anisotropy $A$. From top to bottom, $A = 3.0D^2/J_{ex}$ and $A = 3.75D^2/J_{ex}$. Meron (blue) loops containing an antimeron (red) are clearly visible. Insets show the Fourier transform of $n_z$.

The middle region between two merons with the same chirality. Consequently, a finite scalar spin chirality is accumulated in these regions [see the light green region in Fig. 2(h), Figs. 3(b) and (d)]. The effective interaction between merons decreases with their separation making the meron crystal softer relative to the insertion of topological defects, such as dislocations and disclinations. Indeed, these topological defects are already present in the meron-antimeron crystal shown in Fig. 3(a). Similarly to the case of Abrikosov vortices in superconductors, a small amount of disorder can lead to a meron glass.

The total skyrmion number or scalar spin chirality, $Q_T$, decreases when the easy-plane anisotropy increases. A sufficiently strong easy-plane anisotropy stabilizes a uniformly canted ferromagnetic state with $n_z = B_z/A$. Figure 4 shows a $(B_z, A)$ phase diagram based on the computation of the average skyrmion density.

Our results demonstrate that multiple crystals of mesoscale topological particles can be realized by controlling magnetic anisotropy. The easy-plane anisotropy required to stabilize the meron phase for typical parameters of MnSi, $J_{ex} \approx 3$ meV, $D \approx 0.3$ meV, is $A \approx 0.05$ meV. Our novel meron crystals could be realized in MnSi ultrathin films because the easy-plane anisotropy increases significantly when the film thickness is reduced. Meanwhile the shape anisotropy due to the dipolar interaction also generates an easy-plane anisotropy. Alternatively, spin anisotropy can be controlled by growing the chiral magnet films on heavy element substrates, such as iridium. Note that the meron and antimeron lattice is not connected adiabatically to the skyrmion lattice by tuning the easy plane anisotropy, because the total topological charge is different. We obtain these configurations by annealing. It is difficult to measure the topological charge of the meron directly. One may employ Lorentz transmission electron microscope to visualize the spin orientation directly, which can be used to calculate the topological charge. The skyrmion and meron lattice can also be imaged by the mag-
magnetic force microscope. \cite{29}

To summarize, we have demonstrated how a skyrmion lattice dissociates into a meron and antimeron crystal with increasing easy plane anisotropy in chiral magnets, which can be achieved in MnSi thin films with progressively smaller thickness. The simulated crystal structures are consistent with compact packings of disks of two different radii \cite{25}. Treating topological charge density as an order parameter we can model the triangular to square lattice transition within the Landau theory of phase transitions. We expect a quite different magnetic and transport response from a meron lattice that could usher into novel sensing of very small local magnetic fields as well as information manipulation and storage applications.

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APPENDIX: NUMERICAL METHODS

We employ the Landau-Lifshitz-Gilbert equation of motion for spin $\mathbf{n}$ in the presence of a noisy magnetic field to perform numerical annealing

$$\partial_t \mathbf{n} = -\gamma \mathbf{n} \times \left( \mathbf{B}_{\text{eff}} + \mathbf{B} \right) + \alpha \partial_t \mathbf{n} \times \mathbf{n},$$  \hspace{1cm} (2)

where $\mathbf{B}_{\text{eff}} \equiv -\partial H/\partial \mathbf{n} = J_{\alpha} \nabla^2 \mathbf{n} - 2D \nabla \times \mathbf{n} + \mathbf{B} - A \hat{z}$, $\gamma = 1/(\hbar \gamma)$, and $s$ the magnitude of local spins. Here $\hat{z}$ is a unit vector along the $z$ direction and $\alpha$ is the Gilbert damping coefficient. $\mathbf{B}$ is the fluctuating magnetic field that satisfies the fluctuation-dissipation theorem:  \hspace{1cm} (3)

$$\langle \dot{B}_\mu (\mathbf{r}, t) \dot{B}_\nu (\mathbf{r}', t') \rangle = \frac{2k_B T \gamma}{\gamma^2} \delta_{\mu\nu} \delta(\mathbf{r} - \mathbf{r}') \delta(t - t'),$$

where $T$ is the temperature, $d$ is the film thickness, $k_B$ is the Boltzmann constant and $\mu, \nu = x, y, z$. We have performed numerical annealing starting from high temperatures such that the system is initially in the paramagnetic phase and then goes to zero temperature slowly.

We use dimensionless units in simulations: length is in units of $J_{\alpha} a/D$; energy is in units of $J_{\alpha}^2 D$; magnetic field and easy plane anisotropy are in units of $D^2/J_{\alpha}$; time is in units of $J_{\alpha}/(\gamma D^2)$; temperature is in units of $J_{\alpha}^2/(DK_B)$. A square simulation box with periodic boundary conditions is used. To accommodate the long wavelength structure in Figs. 2 and 3 we have used different system sizes. For the configurations in Fig. 3 we have not reached an ordered state even by changing the system size. The system is discretized with a grid size 0.4$a_0 a/D$ and a smaller grid size is also used to check the accuracy of the results. Equations (2) and (3) are solved using a method introduced in Ref. \cite{30}.
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