

## Antichiral order and spin reorientation transitions of triangle-based antiferromagnets

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We show that triangle-based antiferromagnets with “antichiral” order display a nontrivial dependence of the spin orientation with an in-plane field. The spins evolve from rotating in the opposite sense to the field at very low fields to rotating in the same sense as the field above some critical field scale. In the latter regime the system displays first-order transitions at which the spin angles jump, and these first-order lines terminate in critical points in the Ising universality class. Application to  $\text{Mn}_3\text{Sn}$  is discussed.

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**Introduction.** The elementary unit of a triangle of spins is often considered the building block of frustrated magnetism. Three spins on such a triangle with antiferromagnetic Heisenberg interactions enjoy, in the classical limit, an  $O(3)$  rotational degeneracy of ground states in which the spins lie in a plane at  $120^\circ$  angles to one another. When a field is applied, the degeneracy enlarges from this symmetry-mandated one to an accidental degeneracy which includes both coplanar and noncoplanar states. When such triangles are assembled into the canonical triangular lattice [1], thermal and quantum fluctuations are known to break this degeneracy in favor of the coplanar ones, a phenomenon known as “order by disorder” [2]. Larger degeneracies are found when the triangles are more weakly connected, as in the famous kagome lattice. There, the Heisenberg degeneracy becomes extensive, and ordering is strongly suppressed. Commonly in real materials, weak symmetry-breaking effects such as Dzyaloshinskii-Moriya (DM) coupling [3] and single-ion anisotropy (SIA) provide another degeneracy-breaking mechanism leading to a selection of three-sublattice ordered states.

In this Letter, we study a very common situation of three-sublattice order based on triangles in which the Heisenberg  $O(3)$  symmetry of the Hamiltonian is weakly broken by DM and SIA in favor of coplanar order in zero applied field. We adopt a symmetry approach based on order parameters, which is more universal than microscopic models of specific exchange interactions, but which incorporates a hierarchy of coupling strengths. In particular, we assume that Heisenberg exchange  $J$  is the largest scale, followed by DM with strength  $D$  and SIA of strength  $K$ , i.e.,  $J \gg D \gg K$ . This is inspired by the breathing kagome lattice materials  $\text{Mn}_3\text{Sn}$  and  $\text{Mn}_3\text{Ge}$  [4–10], but is very typical for third row transition-metal magnets. We focus on the *antichiral* state [selected for  $D > 0$ , see Eq. (5)], in which, proceeding clockwise around the triangle, spins rotate *counterclockwise* (Fig. 1). Spins in the antichiral state are nearly free to rotate globally (see below). From this perspective, we consider the evolution of the spin configurations in an applied field, and in particular how the spins rotate when the field is rotated within the XY plane favored by DM

coupling. We show that there is an emergent low magnetic field scale  $H^*$  separating two distinct behaviors. When the field is much smaller than  $H^*$ , the angle of a single spin within the plane rotates in the opposite sense as the field, i.e., if the field is oriented at an angle  $\theta$  in this plane, each spin rotates with angle  $\phi_n = \phi_n^{(0)} - \theta$ , where  $\phi_n^{(0)}$  is an offset for each sublattice  $n$ . Conversely, when the field is much larger than  $H^*$ , the spins rotate in sync with the field, i.e.,  $\phi_n = \phi_n^{(0)} + \theta$ . These competing tendencies result in abrupt and discontinuous changes in the spin configurations, which form lines of first-order transitions in the plane of the magnetic field, terminating at second-order Ising critical points (see Fig. 3). We argue that features recently observed in sensitive measurements of the angular dependence of magnetization and torque in  $\text{Mn}_3\text{Sn}$  [5] are precursors of these transitions, and that the transitions should be observable in higher magnetic fields. The theory should also be applicable to other kagome magnets with antichiral order, of which there are several examples [11–14].

**Symmetry and order parameters.** We begin by presenting a derivation of the free energy as a function of spin angle based on symmetry and the hierarchy of energy scales. We assume at the outset that we have a magnetic system whose ordered structure is fully specified by giving the orientation of a set of three spins on an elementary triangle. We further assume that the dominant interactions between these three spins are antiferromagnetic and isotropic,

$$H_0 = J(\mathbf{S}_0 \cdot \mathbf{S}_1 + \mathbf{S}_1 \cdot \mathbf{S}_2 + \mathbf{S}_2 \cdot \mathbf{S}_0), \quad (1)$$

with  $J > 0$ . This favors an ordered state in which the three spins sum to zero, and have equal magnitudes,

$$\langle \mathbf{S}_n \rangle = \text{Re}[\mathbf{d} e^{\frac{2\pi i n}{3}}], \quad (2)$$

where the angular brackets denote the usual equilibrium thermal expectation value, and  $\mathbf{d}$  is a complex vector satisfying

$$\mathbf{d} \cdot \mathbf{d} = 0. \quad (3)$$

The last relation implies that the magnitude  $|\langle S_n \rangle|$  is independent of  $n = 0, 1, 2$ . This is not a requirement, but would be expected at low temperature classically.

Note that although at a microscopic model level the spins may be taken as either quantum operators or fixed length vectors (classically), the order parameter  $\mathbf{d}$  as defined is always a classical variable, which enables a simple analytic treatment, which is valid regardless of the microscopic formulation.

In general, the order parameter can also be written in terms of two orthogonal real vectors of equal magnitude,  $\mathbf{d} = \mathbf{u} + i\mathbf{v}$ , where Eq. (3) implies  $|\mathbf{u}|^2 = |\mathbf{v}|^2$ ,  $\mathbf{u} \cdot \mathbf{v} = 0$ . These two vectors  $\mathbf{u}$ ,  $\mathbf{v}$  define a plane in which the spins lie. From this one can define a third vector normal to the plane,  $\mathbf{w} = \mathbf{u} \times \mathbf{v} = \frac{1}{2}\text{Im}(\mathbf{d}^* \times \mathbf{d})$ .

In the following, we use the assumption that perturbations from the Heisenberg limit, i.e., DM, SIA, and applied field, are all small compared to  $J$ . Then deviations from the above form are small and more importantly they can be considered to be induced by the perturbations. In the effective field theory (or Landau) sense such deviations correspond to massive modes or subdominant order parameters, and can be integrated out order by order in the perturbations. This allows one to work with an effective free energy which is a function of  $\mathbf{d}$  only [satisfying Eq. (3)], but in which the strength of perturbations may enter nonlinearly.

We construct this free energy based on symmetry and the hierarchy of interactions. First, consider the symmetries in the isotropic limit where  $D = K = 0$ , i.e., with the Hamiltonian in Eq. (1). There is in this case a global  $\text{SO}(3)$  symmetry under  $S_n \rightarrow \mathbf{O}S_n$ , where  $\mathbf{O}$  is an arbitrary  $\text{SO}(3)$  matrix. From Eq. (2), this takes  $\mathbf{d} \rightarrow \mathbf{O}\mathbf{d}$ . Second, Eq. (1) has full  $S_3$  symmetry under arbitrary permutations of the three spins. It is convenient to regard the permutation symmetry as generated by a  $\mathbb{Z}_3$  cyclic permutation which takes  $S_n \rightarrow S_{n+1}$  and a  $\mathbb{Z}_2$  permutation which exchanges  $S_1 \leftrightarrow S_2$ . Under these two operations, respectively, we have  $\mathbf{d} \rightarrow e^{2\pi i/3}\mathbf{d}$  and  $\mathbf{d} \rightarrow \mathbf{d}^*$ .

In zero magnetic field, the only nonzero invariant [using Eq. (3)] under all these symmetries is  $\mathbf{d}^* \cdot \mathbf{d}$ , so the zero-field free energy in the isotropic limit is a function of this quantity alone. This dependence can be regarded as simply fixing the overall magnitude of the order parameter,  $\mathbf{d}^* \cdot \mathbf{d} = 2n_0^2$ , where  $n_0$  is the size of a local moment. While this may shift slightly as anisotropy and field are turned on, the effects can be absorbed in other terms, and we can treat it, following the Landau logic, as fixed.

With this understanding, we now introduce the magnetic field  $\mathbf{h}$  on the isotropic spins. It transforms in the same way as  $\mathbf{d}$  under global  $\text{SO}(3)$  rotations, and is invariant under all the permutations. Consequently, we find that the purely field-induced terms in the free energy are of the form

$$F_h^{\text{iso}} = c_1 |\mathbf{h} \cdot \mathbf{d}|^2 + c_2 \text{Re}[(\mathbf{h} \cdot \mathbf{d})^3] + O(h^4). \quad (4)$$

As is typical for an antiferromagnet, there is no linear coupling of the field to the order parameter, but in this case both quadratic and cubic terms occur [15, 16].

The physical meaning of these terms is as follows. The leading quadratic term selects configurations in which the spins lie in a plane either normal to or containing the field, for  $c_1 > 0$  and  $c_1 < 0$ , respectively. Note that the form in Eq. (2) only defines the antiferromagnetic components of the

spins (the primary order parameter), and not the field-induced uniform moment. For the semiclassical Heisenberg antiferromagnet on the triangular lattice, the two types of orderings are classically degenerate (i.e., at  $1/S = T = 0$ ), but it is known that the coplanar configurations are favored by both thermal and quantum fluctuations, which select  $c_1 < 0$  [ $c_1 \sim -1/(JS)$  at  $T = 0$ ] [17]. The cubic term selects an orientation of the spins within this plane: When the sign of  $c_2$  is positive (negative), one of the three spins lies antiparallel (parallel) to the field. According to Ref. [17], for the triangular lattice the preferred configuration of the former type, and  $c_2 > 0$  [ $c_2 \sim 1/(J^2S)$  in at  $T = 0$ ]. The same signs are found for the classical kagome lattice at nonzero temperature due to thermal fluctuations (though the estimates differ quantitatively due to the higher degeneracy of the kagome case) [18].

Now consider the effects of DM and SIA, of the microscopic form

$$H' = \sum_n [D\hat{\mathbf{z}} \cdot \mathbf{S}_n \times \mathbf{S}_{n+1} - K(\hat{\mathbf{e}}_n \cdot \mathbf{S}_n)^2], \quad (5)$$

where  $\hat{\mathbf{e}}_n = [\cos(\frac{2\pi n}{3}), \sin(\frac{2\pi n}{3}), 0]$ . These additional terms lower the symmetry as follows. The DM interaction  $D$  maintains a global  $\text{SO}(2)/\text{U}(1)$  subgroup of  $\text{SO}(3)$  under rotations about the  $\hat{\mathbf{z}}$  axis, under the  $\mathbb{Z}_3$  cyclic permutation of the spins, and under the spin-orbit coupled  $C_2$  symmetry in which the  $\mathbb{Z}_2$  spin permutation discussed earlier is combined with the corresponding rotation in spin space,

$$C_2 : S_0 \rightarrow O_2 S_0, \quad S_{1/2} \rightarrow O_2 S_{2/1}, \quad (6)$$

where  $O_2 = \text{diag}(1, -1, -1)$ . With the SIA term  $K$ , the symmetry is further reduced, so that the global  $\text{SO}(2)$  and  $\mathbb{Z}_3$  operations are collapsed to a single  $C_3$  combined rotation

$$C_3 : S_n \rightarrow O_3 S_{n+1}, \quad (7)$$

where  $O_3$  is the appropriate rotation matrix [19].

To incorporate the symmetry-lowering effects, it is convenient to adopt a new basis

$$d_{\pm} = \frac{1}{2}(d_x \pm id_y), \quad (8)$$

and trade  $\mathbf{d}$  for  $d_+$ ,  $d_-$ , and  $d_z$ . Note that because  $\mathbf{d}$  is complex,  $d_+$  is *not* the conjugate of  $d_-$  and is an independent complex field. The symmetry operations in the new basis become

$$\begin{aligned} \text{SO}(2) : d_+ &\rightarrow e^{i\theta} d_+, d_- \rightarrow e^{-i\theta} d_-, d_z \rightarrow d_z, \\ \mathbb{Z}_3 : d_+ &\rightarrow e^{2\pi i/3} d_+, d_- \rightarrow e^{2\pi i/3} d_-, d_z \rightarrow e^{2\pi i/3} d_z, \\ \mathbb{Z}_2 : d_+ &\rightarrow d_+^*, d_- \rightarrow d_-^*, d_z \rightarrow d_z^*, \\ C_3 : d_+ &\rightarrow e^{4\pi i/3} d_+, d_- \rightarrow d_-, d_z \rightarrow e^{2\pi i/3} d_z, \\ C_2 : d_+ &\rightarrow d_+^*, d_- \rightarrow d_-^*, d_z \rightarrow -d_z^*, \\ \mathcal{T} : d_+ &\rightarrow -d_+, d_- \rightarrow -d_-, d_z \rightarrow -d_z. \end{aligned} \quad (9)$$

To summarize, the DM term is invariant under  $\text{SO}(2)$ ,  $\mathbb{Z}_3$ ,  $C_2$ , and  $\mathcal{T}$ . The  $K$  term is invariant under  $C_3$ ,  $C_2$ , and  $\mathcal{T}$ . So the first three lines above are approximate symmetries while the final three are exact. It is also useful to note that under  $\mathbb{Z}_2$ ,  $D \rightarrow -D$  (but  $\mathbb{Z}_2$  does not act simply upon  $K$ ).

Using the above symmetries, and using the constraint Eq. (3) and the condition that  $\mathbf{d}^* \cdot \mathbf{d} = 2n_0^2$ , the most general,

nonconstant free-energy terms at zero field and quadratic in  $\mathbf{d}$  are

$$F_2 = s_1(d_+^*d_+ - d_-^*d_-) + s_2(d_z^*d_z - 2d_+^*d_+ - 2d_-^*d_-) + s_3 \text{Re}(d_-^2). \quad (10)$$

A naive calculation, simply inserting Eq. (2) in Eq. (5), shows that  $s_1 \sim -D$ ,  $s_3 \sim -K$ , while  $s_2 \sim K$ , while an additional contribution  $\Delta s_2 \sim D^2/J$  is expected to arise at second order in the DM coupling (it must be even in  $D$  because under  $\mathbb{Z}_2 D \rightarrow -D$  but  $s_2$  is invariant). In both cases  $s_2, s_3 \ll |s_1|$ .

Here, we are interested in  $D > 0$  which implies  $s_1 < 0$  which favors  $d_- = d_z = 0$  and  $|d_+| = n_0$ . This is the antichiral state. Then the  $s_2$  term is constant and the  $s_3$  term vanishes. Note that the phase of  $d_+$  is arbitrary at this level, reflecting the fact that the rotation of the spins is in the opposite sense to the rotation of the local easy axes, so that the two are incompatible. If by contrast we take  $D < 0$ , the chiral state with  $d_+ = d_z = 0$  is stabilized and  $|d_-| = n_0$ . Then the  $s_3$  term is nonzero and in fact fixes the phase of  $d_-$ , which means the spins are not free to rotate in the chiral state.

For the antichiral state, the complete freedom to rotate the phase is an artifact of the truncation of Eq. (10) to second order in  $\mathbf{d}$ . A nontrivial invariant fixing the phase of  $d_+$  arises at sixth order:

$$f_6 = \lambda \text{Re}(d_+^6). \quad (11)$$

We expect that  $\lambda \sim K^3/J$  (because it is sixth order in  $d_+$ , and hence in spin expectation values, while the  $K$  term is quadratic in spins) as was verified by calculations for  $\text{Mn}_3\text{Sn}$  [20], and therefore is extremely small and often negligible.

Now consider the terms involving the magnetic field. Similarly to Eq. (8), define

$$h_{\pm} = h_x \pm ih_y. \quad (12)$$

Note that  $h_{\pm}^* = h_{\mp}$  (unlike for  $d_{\pm}$ ), so it is sufficient to list the properties of  $h_+$  and  $h_z$ . Under the various transformations, we have

$$C_3 : h_+ \rightarrow e^{2\pi i/3}h_+, \quad h_z \rightarrow h_z, \quad (13)$$

$$C_2 : h_+ \rightarrow h_+^* = h_-, \quad h_z \rightarrow -h_z, \quad (14)$$

$$\mathcal{T} : h_+ \rightarrow -h_+, \quad h_z \rightarrow -h_z. \quad (15)$$

Comparing now Eqs. (9) and (13), we can find invariants involving the field and  $d_{\mu}$ . To linear order in the field, we find

$$f_{h,1} = g_1 \text{Re}(h_+d_+) + g_2 h_z \text{Im}(d_-). \quad (16)$$

In  $\text{Mn}_3\text{Sn}$ , where the order is antichiral,  $d_- = 0$  and only the  $g_1$  term is active. It is order of  $g_1 \sim K/J$  [20], because it requires nonzero  $K$  to break the  $\text{SO}(2)$  symmetry. We see that the linear coupling to the field multiplies  $h_+$  and  $d_+$ , which favors rotating these complex numbers *with opposite phases*. This expresses the surprising phenomenon that each spin in the antichiral case at small fields actually rotates in the opposite sense as the applied field. The  $g_1$  term can be understood physically from the picture in Fig. 1: The net effect of a rigid counterclockwise  $C_3$  rotation of the antichiral triangle is the same as rotating the spins in place by a clockwise  $C_3$  rotation.

Note that this effect contradicts the behavior in the isotropic system, which is dictated by the cubic coupling in

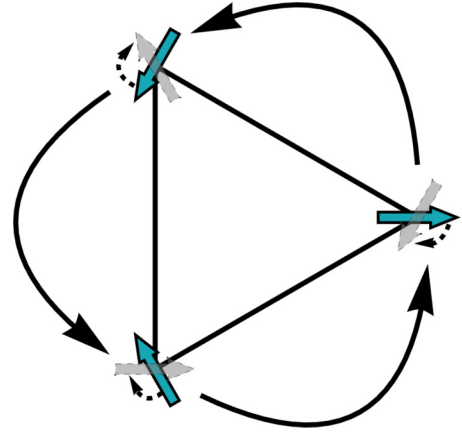


FIG. 1. Illustration of the behavior of antichiral spins under rotations. A *counterclockwise* rigid rotation of the triangle by  $120^\circ$ , indicated by the solid arrows, rotates the spins from their initial state (shown in blue) to a new state, shown in gray. This is equivalent to a *clockwise* rotation of the spins by  $120^\circ$  (indicated by the smaller dashed arrows) in place.

Eq. (4), and favors rotating each spin in sync with the field. The opposite tendencies lead to a transition as a function of field strength.

To unveil it more cleanly, we focus now on the case in which only  $d_+$  is assumed nonzero, and the magnetic field is in the plane, and write the free energy as a series in  $d_+$  and the field only. We furthermore assume that the higher-order terms in the field are dominated by the ones already present in Eq. (4), and simply express those in the case where  $d_z = d_- = 0$  in terms of  $d_+$ . We find in this case  $\mathbf{h} \cdot \mathbf{d} = 1/2 h_- d_+$  which leads to

$$f_+ = \lambda \text{Re}(d_+^6) + g_1 \text{Re}(h_+d_+) + \frac{c_1}{4} |h_+|^2 |d_+|^2 + \frac{c_2}{8} \text{Re}[(h_-d_+)^3]. \quad (17)$$

Equation (17) is *not* the most general free energy allowed by symmetry, but rather the minimal one which includes the largest terms which break all approximate symmetries both at zero and nonzero fields. In particular,  $\lambda$  and  $g_1$  are the leading order (in  $K$ ) terms that break  $\text{SO}(2)$  symmetry at zero field and nonzero fields, respectively, as seen from, e.g., Ref. [20].

To analyze Eq. (17), we change to angular coordinates,  $h_+ = h e^{i\theta}$  and  $d_+ = d e^{i\phi}$ . It becomes, up to a constant,

$$f_+ = -w \cos 6\phi - uh \cos(\phi + \theta) - vh^3 \cos 3(\phi - \theta), \quad (18)$$

where

$$w = -\lambda d^6, \quad u = -g_1 d, \quad v = -c_2 d^3/8. \quad (19)$$

*Angular analysis and phase transitions.* Equation (18) is the general result for the angle-dependent free energy of the antichiral state. We now show that it exhibits the phase transitions described in the Introduction. Without loss of generality, we take  $u, v, w > 0$ . Using the aforementioned estimates  $w \sim K^3/J^2$ ,  $u \sim K/J$ , and  $v \sim D/J^3$ , we establish the condition  $w \ll \sqrt{u^3/v}$ , under which  $w$  can be neglected in the field regime  $h \gg w/u$ . We henceforth assume this condition and take  $w = 0$ . Then the order parameter angle  $\phi$  is determined

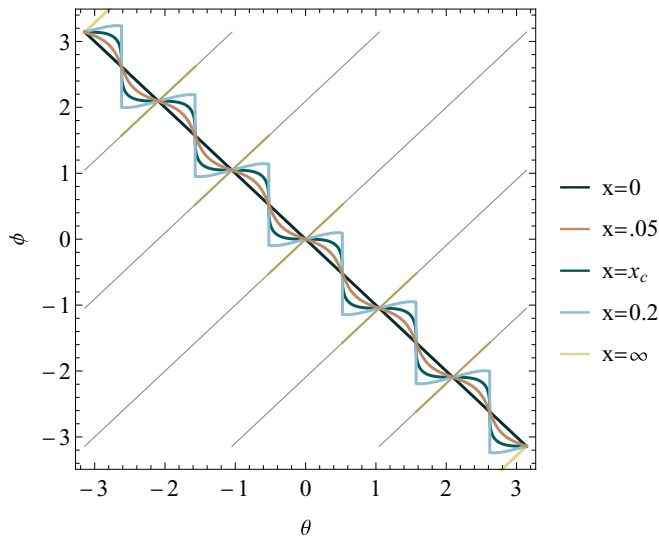


FIG. 2. Spin angle  $\phi$  vs field angle  $\theta$  for various values of  $x$ . The thin gray lines show the three degenerate branches at  $x = \infty$ , while the piecewise linear function labeled  $x = \infty$  is the one selected at asymptotically large but finite  $x$ .

just by minimizing the final two terms in Eq. (18). So we may write  $f_+ = uhg(\phi, \theta)$ , with

$$g(\phi, \theta) = -\cos(\phi + \theta) - x \cos 3(\phi - \theta), \quad (20)$$

where  $x = \sqrt{v/uh^2} > 0$ . The optimal spin angle  $\phi(\theta, x)$  is determined from minimizing  $g$  at fixed field angle  $\theta$  and  $x$ . It is instructive to analyze the two limits  $x = 0$  and  $x = \infty$ . At  $x = 0$ ,  $g$  is clearly minimized by  $\phi = -\theta$ . At  $x = \infty$ , there are three degenerate minima with  $\phi = \theta + 2\pi k/3$ , with  $k = 0, 1, 2$ . One observes that the spin angles wind in the opposite sense in the two extreme limits. The degeneracy in the large  $x$  limit is resolved by selecting the branch ( $k$ ) which minimizes the first term in Eq. (20). This leads to jumps in  $k$  (and hence  $\phi$ ) as a function of  $\theta$ , which occur when  $\theta = \pi/6 + \pi m/3$ , with integer  $m$ , as shown in Fig. 2. Away from the two extreme limits, the curve  $\phi(\theta)$  evolves, but the discontinuities persist for large  $x$ , while they are absent for small  $x$ . A transition occurs for  $x = x_c$ , where the discontinuities first appear.

To clarify the critical points, we define  $\psi = \phi + \theta$ , and let  $\theta = \pi/6 + \omega$ , so that  $\omega = 0$  defines the location of one of the discontinuities for  $x > x_c$ . Some algebra gives

$$\begin{aligned} \tilde{g}(\psi, \omega) &= g(\psi - \frac{\pi}{6} - \omega, \theta) \\ &= -\cos \psi + x \cos 6\omega \cos 3\psi + x \sin 6\omega \sin 3\psi. \end{aligned} \quad (21)$$

We see that at  $\omega = 0$ ,  $\tilde{g}(\psi, 0)$  is an even function of  $\psi$ . In fact, the evenness of this function reflects a symmetry  $\phi \rightarrow \pi/3 - \phi$ , which is a  $C_2$  symmetry of the Hamiltonian when the field angle  $\theta = \pi/6$ . Hence  $\tilde{g}(\psi, \omega)$  can be regarded as a Landau function, with a minima at  $\psi = 0$  for  $x < x_c = 1/9$ , which bifurcates for  $x > x_c$  into two degenerate minima at  $\psi = \pm\psi_0$ . This is an Ising phase transition. Note that for  $\text{Mn}_3\text{Sn}$ , the Ising symmetry along these lines can be traced to a symmetry under a  $C_2$  rotation about the axis of the magnetic field, when the field is aligned to these crystalline axes. The deviation  $\omega$  of the field angle plays the role of a

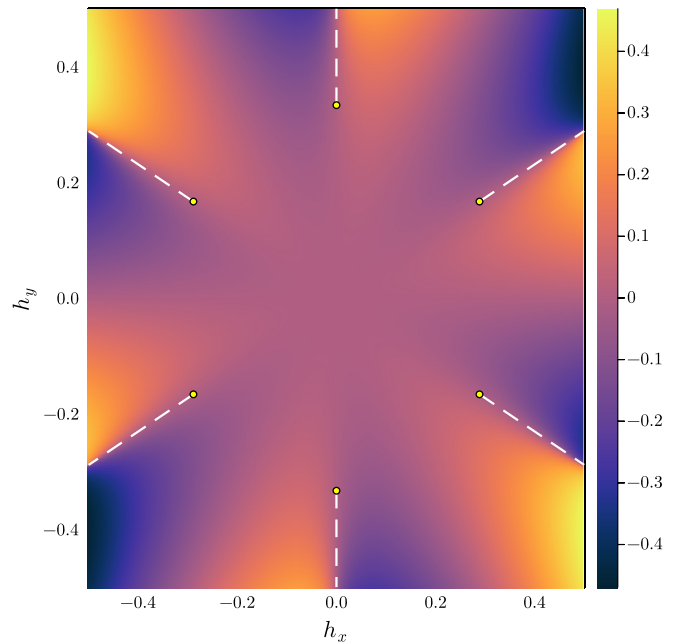


FIG. 3. Density plot of torque ( $\tau/u$ ) vs in-plane magnetic field. The field scale is chosen so that  $h_c = 1/3$ . Yellow circles mark the location of Ising critical points, and the first-order lines are marked with dashed lines.

symmetry-breaking field on the Ising order parameter, and the discontinuities in  $\theta$  are analogous to the first-order transition that occurs within the ordered phase of the Ising model on changing the sign of the field. This analysis determines a critical field  $h_c = \frac{1}{3}\sqrt{\frac{\pi}{v}}$ . Following the analogy with the Ising model, one observes within this mean field picture that at the critical field, the angle  $\psi \sim -|\omega|^{1/3} \text{sgn}(\omega)$  for small variations of the field angle near  $\pi/6$ . This is analogous to the nonlinear susceptibility of the Ising ferromagnet at criticality. *A priori* thermal fluctuations will renormalize this exponent to that of the three-dimensional (3D) Ising model—there we have  $\psi \sim |\omega|^{1/\delta}$  and the true value is approximately  $\delta \approx 4.8$ .

One can readily calculate the torque  $\tau = df/d\theta$ , where the *total* derivative should be taken with respect to the field angle of the energy optimized at  $\phi = \phi(\theta)$ . Using the energy minimization condition, one obtains  $\tau = 2uh \sin[\phi(\theta) + \theta]$ , which is plotted in Fig. 3. It is discontinuous across the first-order transition lines, and vanishes when  $\theta$  is a multiple of  $\pi/3$ .

The angular transition may also be detectable by means other than the torque. Hysteresis may occur and domains may form near the first-order lines. The spin angles themselves could be observed directly in neutron scattering. It is interesting to point out that a very similar phenomenon has already been observed in  $\text{CeAlGe}$ , in which the spontaneous formation of domains was argued to give rise to a sharp peak in the resistance versus angle, dubbed singular angular magnetoresistance [21]. It would be very interesting to study the angular dependence of the resistance in  $\text{Mn}_3\text{Sn}$  in an appropriate range.

From the calculations in Ref. [20], and the Supplemental Material of Ref. [5], we can compare directly to Eq. (18), and



extract the parameters of the symmetry-based theory in terms of microscopics. In the classical zero-temperature model, one obtains thereby

$$u = \frac{K}{J + \sqrt{3}D}, \quad v = \frac{D}{3\sqrt{3}(J + \sqrt{3}D)^3}. \quad (22)$$

This leads to the critical field, restoring units

$$H_c = \frac{J + \sqrt{3}D}{g\mu_B} \sqrt{\frac{K}{D}}. \quad (23)$$

Taking  $D = 0.2J$ ,  $K = 0.006J$ ,  $J = 20$  meV, and  $g = 3$  yields  $H_c \approx 20T$ . This is of course to be renormalized by thermal and quantum fluctuations, but gives an idea of the order of magnitude. It strongly suggests the transition should be within the range of current experiments.

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