Chiral multiple-q and skyrmion phases induced by Rashba-Hund interactions in the kagome lattice

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Engineering non trivial topological phases in materials, specially skyrmion-like arrangements, has been of great interest in the last decade due to its potential technological applications. In this work, we study a model of electrons coupled to a magnetic texture in the kagome lattice interacting with magnetic moments via Rashba spin orbit coupling in the large Hund interaction limit. We obtain the effective spin Hamiltonian and study the emergent low temperature phases under an external magnetic field using large scale Monte Carlo simulations. We show that strong geometric frustration, characteristic of the kagome lattice, and the competition between the effective exchange, antisymmetric and anisotropic couplings, gives rise to a large variety of non-trivial topological phases. On the one hand, for antiferromagnetic exchange coupling a pseudo-antiferromagnetic skyrmion crystal is stabilized for a broad range of parameters. As the exchange coupling gets smaller, chiral single-q and double-q phases emerge. On the other hand, in the ferromagnetic case, even though the competition of the different interactions produce a series of exotic textures, a remarkable parallel may be drawn with the pure ferromagnetic model with antisymmetric interactions. Finally, for the special case where the exchange coupling is completely suppressed, we show that, coming from a higher temperature cooperative paramagnet, an “umbrella-like” plaquette order with semiextensive degeneracy is induced by the external magnetic field.

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I. INTRODUCTION

Undoubtedly, the search for new materials and their application to technological devices is a signature of this information driven era. One of the new exciting proposals is the possibility of using magnetic skyrmions in race track memory devices. Given that skyrmions can be effectively moved with electrical currents but are deflected from their trajectory due to the “skyrmion Hall effect”, an extensive amount of attention has gone into the study of antiferromagnetic skyrmions, where this effect should be suppressed. Experimentally, the inhibition of this Hall effect has been observed for ferromagnetic skyrmions in GdFeCo films, synthetic antiferromagnetic skyrmion bubbles, and it is also expected to be absent for other exotic structures such as skyrmionium.

The quest for antiferromagnetic skyrmion-like structures has certainly been yielding promising results. This type of topological structures have been recently detected in different materials; for example, they have been stabilized at room temperature in synthetic antiferromagnets, antiferromagnetic half-skynrmions and skyrmions and bimerons have been realized in \( \alpha \)-Fe$_2$O$_3$; fractional antiferromagnetic skyrmion crystals have been shown to be stabilized in MnSc$_2$O$_4$.

Theoretically, multiple models have been proposed. Closely related to this work, a key ingredient in some cases has been an underlying geometrically frustrated lattice. Antiferromagnetic skyrmion crystals formed by interpenetrated skyrmion sublattices have been found to be stable under a magnetic field for different models in the antiferromagnetic triangular lattice, combining Dzyaloshinskii-Moriya interactions and even additional competing terms. Other competing interactions in the frustrated triangular lattice have been shown to produce exotic states such as spontaneous skyrmion lattices and higher order skyrmions.

Quite recently, different types of antiferromagnetic skyrmion crystals have been reported for microscopic electronic models with Rashba spin-orbit coupling (SOC) both in the square and triangular lattices. Rashba SOC had already been shown to be a stabilization mechanism for skyrmions in two dimensional lattices.

On the other hand, we have also previously studied the formation of topologically non trivial structures in a simple model in the antiferromagnetic kagome lattice with Dzyaloshinskii-Moriya (DM) interactions, where we found that the strong frustration in kagome induces a topological phase characterized by a periodical arrangement of three interpenetrated non-trivial sublattices formed by pseudo-skyrmion magnetic textures with topological charge \( |Q| < 1 \).

In this work, our main goal is to explore the effect of the strong geometrical frustration characteristic of the kagome lattice in the possible topological phases that
may emerge at low temperature in Rashba SOC models. Experiments have shown the realization of skyrmions and competing magnetic orders in the centrosymmetric material Gd₃Ru₄Al₁₋ₓGaₓ where the magnetic moments form a breathing kagomé lattice and models combining frustrating exchange couplings and other types of interaction have been proposed.

Motivated by this, here we focus on the emergent low temperature phases induced by an external magnetic field in the spin effective models combining Rashba SOC and strong Hund’s coupling. These effective models have one very interesting feature: besides the well known exchange strong Hund’s coupling, these effective models have one temperature phases induced by an external magnetic field have been proposed.

II. MODEL

We start with the ferromagnetic Kondo lattice model in the presence of Rashba SOC on a kagome lattice described by the following Hamiltonian,

\[
\mathcal{H} = -t \sum_{\langle ij \rangle, \sigma} (\hat{c}_{i\sigma}^\dagger \hat{c}_{j\sigma} + H.c) - i\lambda \sum_{\langle ij \rangle, \sigma, \sigma'} \hat{c}_{i\sigma}^\dagger (\hat{\sigma} \times \hat{e}_{ij})^z_{\sigma\sigma'} \hat{c}_{j\sigma'}
\]

\[- \frac{J_H}{2} \sum_{i} \mathbf{S}_i \cdot (\hat{c}_{i,a}^\dagger \hat{\sigma}_{ab} \hat{c}_{i,b}) \]

where \(\{\hat{c}_{i\sigma}^\dagger, \hat{c}_{i\sigma}\}\) are the fermionic creation and annihilation operators with spin \(\sigma = \pm\) at the site \(i\), \(t\) is the nearest neighbor hopping amplitude, \(\lambda\) is the strength of the Rashba interaction, \(\hat{e}_{ij}\) are the unitary vectors pointing from site \(i\) to \(j\) in the kagome lattice (see Fig. 1) and \(\hat{\sigma}\) represents the Pauli matrices vector. \(J_H\) denotes the Hund’s coupling parameter and \(\mathbf{S}_i\) is the classical localised spin at site \(i\).

Assuming \(J_H \gg t\) and taking the double-exchange approximation at half-filling, we derive, via second order perturbation theory\([23]\), the following effective spin model

\[
\mathcal{H}_{\text{eff}} = -\frac{1}{J_H} \sum_{ij} [t^2(1 - \mathbf{S}_i \cdot \mathbf{S}_j) - 2t\lambda \hat{\gamma}_{ij} (\mathbf{S}_i \times \mathbf{S}_j)]
\]

\[+ \lambda^2 (1 + \mathbf{S}_i \cdot \mathbf{S}_j - 2(\hat{\gamma}_{ij} \mathbf{S}_i) (\hat{\gamma}_{ij} \mathbf{S}_j)) \]

(2)

where \(\hat{\gamma}_{ij} = \hat{z} \times \hat{e}_{ij}\) are unitary vectors and \(\mathbf{S}_i\) is the local three component magnetic moment at site \(i\).

In order to study the competition between the hopping and Rashba parameters in this effective model, it is convenient to describe all the interactions by one parameter, \(\alpha\), such that the hopping and Rashba couplings are rewritten as \(t = (1 - \alpha)t_0\) and \(\lambda = \alpha t_0\) (\(t_0 = 1\) sets the reference energy scale). Now, we define the exchange coupling \(J = 1 - 2\alpha\), the DM interaction \(D = D\hat{\gamma}_{ij}\) with \(D = 2\alpha(1 - \alpha)\) (\(D = 2\alpha\)), and a bond anisotropy term \(J_{\parallel} = 2\alpha^2\) (\(J_{\perp} = 2\lambda\)). The effective model including the presence of an external magnetic field of strength \(h_z\) in the \(\hat{z}\) direction reads as

\[
\mathcal{H}_{\text{eff}} = \frac{t^2}{J_H} \sum_{ij} [J \mathbf{S}_i \cdot \mathbf{S}_j + D \hat{\gamma}_{ij} \cdot (\mathbf{S}_i \times \mathbf{S}_j)]
\]

\[+ J_{\parallel} (\hat{\gamma}_{ij} \cdot \mathbf{S}_i) (\hat{\gamma}_{ij} \cdot \mathbf{S}_j)] - h_z \sum_{i} S^z_i \]

(3)

In Fig. 2 we show the evolution of the effective interactions as a function of \(\alpha\), where it can be appreciated how the model is predominantly antiferromagnetic for \(\alpha < 0.5\) (dominated by \(J > 0\)) and it becomes ferromagnetic with large bond anisotropy \(J_{\parallel}\) for \(\alpha > 0.5\). The special point \(\alpha = 0.5\) corresponds to the value where the exchange coupling \(J = 0\), the DM interaction takes its maximum value and the bond anisotropy satisfies \(J_{\parallel} = D\).

In the following section, we explore in detail these different regimes of the interactions and the low temperature emergent phases using large-scale Monte Carlo simulations. For simplicity, we fix \(t_0^2/J_H = 1\) and take \(|J| \equiv 1\) to set the temperature scale, except in the special case \(\alpha = 0.5\), \(J = 0\). In the following, we describe the numerical simulations and the relevant parameters that we have used to build the different phase diagrams as a function of magnetic field \(h_z\) and \(\alpha\).

III. MONTECARLO SIMULATIONS

To explore the low temperature phases of the models presented in Eq. (3), we perform Monte Carlo simulations with the Metropolis algorithm using a combination of overrelaxation and heat bath methods in an annealing scheme. We run our simulations in \(3 \times L^2\) site clusters, \(L = 24 - 48\), with periodic boundary conditions. \(10^5 - 10^6\) Monte Carlo steps (MCS) were used for initial relaxation, and measurements were taken in twice as much MCS, scanning each \(\alpha\) value in a large region of external magnetic fields.

In order to characterize the possible magnetically ordered states of model in Eq. (3), we have calculated various physical observables to precisely identify the different phases. First, we have computed the static spin structure factor with components perpendicular \(S^z_q\) to the external...
FIG. 1: Kagome lattice. Labels $i,j,k$ denote the three triangular sublattices. Orange arrows, perpendicular to the bonds, represent the Dzyaloshinskii-Moriya vectors.

FIG. 2: Effective interactions as a function of $\alpha$: exchange coupling (blue) $J = 1 - 2\alpha$, exchange anisotropy $J_\parallel = 2\alpha^2$, and DM interaction with strength $D = 2\alpha(1-\alpha)$. The peak location of the structure factor contains information about the ordering wave vector. In particular, it is well known that a possible skyrmion phase is characterized by a triple-$q$ structure factor, with six Bragg peaks and $\sum q = 0$, whereas helical or conical phases have single- or double-$q$. In addition, we compute the scalar chirality $\chi$ which is directly linked to the topological charge or skyrmion number. In this work, we have calculated two types of scalar chirality: a nearest-neighbor scalar chirality $\chi_{nn1}$, taken as the sum of the triple product of spins in a triangular plaquette (Eq. (5)), and the third nearest-neighbor scalar chirality $\chi_{nn3}$, computed in each triangular sub-

$$S_q^\perp = \frac{1}{N}(|\sum_i S_i^x e^{i\mathbf{r}_{ij} \cdot \mathbf{r}_i}|^2 + |\sum_i S_i^y e^{i\mathbf{r}_{ij} \cdot \mathbf{r}_i}|^2) \quad (4)$$

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$$\chi_{nn1} = \frac{1}{8\pi} \sum_{m=1}^{N/3} \chi_{mpq} \quad (5)$$

$$\chi_{nn3} = \frac{1}{8\pi} \sum_{a=1}^3 \sum_{m=1}^{N/3} \chi_{(a)mpq} \quad (6)$$

where in Eq. (5) $\chi_{mpq} = \mathbf{S}_m \cdot (\mathbf{S}_q \times \mathbf{S}_p)$ is the local chirality in the triangular plaquettes of the kagome lattice (for example, in an upward triangular plaquette, $m,p,q \equiv i,j,k$ in Fig. 1). In Eq. (6), $\chi_{(a)mpq}$ is the local sublattice chirality, taking the three spins $m,p,q$ in elementary triangles in the three triangular sublattices of the kagome lattice, $a = 1,2,3$. 

FIG. 3: (A) $h_z$ vs $\alpha$ phase diagram for the antiferromagnetic region ($\alpha < 0.5$), built from structure factor calculations at $T \approx 4 \times 10^{-3}$. (B-E) typical structure factors for each phase: locally chiral (B) at $\alpha = 0.1$, $h_z = 1.5$, pseudo-skyrmion at $\alpha = 0.1$, $h_z = 2.5$ (C), double-$q$ at $\alpha = 0.25$, $h_z = 2.5$ (D) and single-$q$ phase at $\alpha = 0.4$, $h_z = 1.5$ (E).
A. Antiferromagnetic Phase Diagram - $\alpha < 0.5 - J > 0$

We first focus on the $\alpha < 0.5$ region, where the exchange coupling is antiferromagnetic, and the DM and $J_\parallel$ interaction increase with $\alpha$. $J$ dominates until $\alpha = -1/2 + \sqrt{3}/2$, where $J = J_\parallel$. Therefore, for lower values of $\alpha$ we may think of the models as antiferromagnetic exchange (AF) models with perturbative DM and bond anisotropy interactions. We have previously studied the case of the pure AF+DM model\textsuperscript{23} (taking in this case DM interactions along the bonds), where the magnetic field $h_z$ induced an antiferromagnetic pseudo-skyrmion crystal (pSkX) phase for a broad region in parameter space. Now, we extend this study by the inclusion of bond anisotropy, analyzing in a way its effect in the pSkX formation.

The main results are presented in Fig. 4, where we show the $h_z$ vs. $\alpha$ phase diagram in panel (A), and the representative structure factor for each phase in panels (B-E). We found that, in addition to the single-q states, various types of multiple-q states are also stabilized under magnetic fields due to the competition between the three types of interactions. For $\alpha \leq 0.2$ and low magnetic fields $h_z < 2$, the DM interaction stabilizes a locally chiral phase presenting six bright peaks around every high-symmetry $M$ point in the spin structure factor (as is shown in Fig. 3B). The non periodic structure of this phase (shown in Fig. 3A) arises because neither $D$ nor $J_\parallel$ are large enough to lift the degeneracy at low magnetic fields.

For larger magnetic fields, $h_z > 2$, an antiferromagnetic pseudo-Skyrmion crystal phase (which we call pSkx-I) emerges, which is recognizable due to the presence of 12 peaks in $S_3$ around the $M$ points in the extended Brillouin zone (EBZ) (panel C) and a spin texture formed by a particularly intricate three-sublattice structure (Fig. 4B). In addition, we found a single-q phase for $\alpha > 0.25$ and low magnetic field ($h_z < 3.5$). Finally, a double-q phase at low $h_z$ emerges for a narrow region of $\alpha$, which broadens as the external field is increased. In the following subsections, we proceed to give a more detailed description of all the low temperature phases in detail.

1. Néel type pseudo-skyrmions

For $\alpha < 0.2$ an antiferromagnetic pseudo-skyrmion crystal (pSkX-I) phase emerges in an intermediate range of magnetic fields. This spin texture is formed by a particular three-sublattice split puzzle, whose structure we studied in detail in a previous work\textsuperscript{23}. The “pseudo-skyrmion” name is given due to the similarity between these structures and skyrmions. However, in this case the center is not fully polarized, and as a consequence of this the topological charges are $|Q| \leq 1$. According to the Hamiltonian in Eq. (5), since the DM interaction is perpendicular to the bonds, the pSkX-I phase hosts Néel type structures. The pSkX-I is formed by three interpenetrated no trivial structure\textsuperscript{22}, and its hidden topological order may be inspected with the third nearest neighbor chirality $\chi_{nn3}$ defined in Eq. (6).

In Fig. 5, we plot $\chi_{nn3}/L^2$ as a function of $h_z$ for $T = 4 \times 10^{-3}$ and $L = 24$ for $\alpha = 0.1, 0.15, 0.2$. As the magnetic field is increased, $\chi_{nn3}$ increases almost linearly with $h_z$ due to the chiral nature of the magnetic domains. In the pSkX-I region, $\chi_{nn3}$ is almost constant, indicating a fixed number of pseudo-skyrmions. Finally, it drops to zero for larger $h_z$ in the double-q phase. Since the skyrmion size depends on the DM strength $D$, it is expected that $\chi_{nn3}$ will change for different values of $\alpha$.

2. Double-q and single-q phases

The double-q phase, that covers a large part of the $h_z$-$\alpha$ phase diagram, presents a structure factor $S_3^\parallel$ with strong intensity in two incommensurate wave vectors $q^\parallel$ which lie between the $K$ points from the Brillouin Zone and the EBZ (Fig. 3D). The single-q phase, which emerges for $\alpha \gtrsim 0.25$, is characterized by one incommensurate wave vector $q^\parallel$ as is shown in Fig. 3C. In general, in most systems hosting skyrmion lattices, the single-q phase corresponds to a helical phase while for the double-
q phase could be simply the superposition of two spirals or helices.

With the aim to characterize these phases, let us begin with the inspection of the corresponding snapshots. On one hand, in the double-q phase, whose spin texture is shown in Fig. 6A, we observe that the spin configurations are composed by non-coplanar “Y-type” structures, alternating with higher-magnetization hexagons. Its real-space sublattice spin configurations present non-equivalent stripes (Fig. 6, panels B-C-D) indicating a sublattice symmetry breaking. Clearly, which sublattice is arranged in which way depends on the MC realization and the lattice symmetry would be restored when averaging on several realizations, as we have checked.

On the other hand, in the single-q phase (a representative spin texture is shown in Fig. 7A), the arrangement is also not symmetric: in two of the sublattices (Fig. 7C-D), the components are homogeneous, while in the remaining sublattice, there is a Néel modulation. Splitting into the three sublattices (panels B-C-D in Fig. 7) it can be seen that these are not equivalent. Then, as a first point, we see that both single-q and double-q phases have broken sublattice symmetry.

Now, to complete the characterization of this exotic single-q and double-q phases, we resort again to the nearest neighbors scalar chirality \( \chi_{nn1} \) (Eq. (5)). As an example, we show \( \chi_{nn1} \) as a function of \( h_z \) for \( T = 4 \times 10^{-3} \), for \( \alpha = 0.35, 0.4, 0.45 \), where the system goes from a single-q to a double-q phase as the field increases. This change is reflected in the \( \chi_{nn1} \) curves: at low fields, \( \chi_{nn1} = 0 \), corresponding to the single-q phase, and as the system is further magnetized, it enters in the double-q phase with a non-zero net chirality. Although the single-q phase has no net chirality, a closer look reveals that the spin arrangement has an alternate chirality, similar to what has been seen in kagome for example in the cuboc phases. This is illustrated in Fig. 7, where it can be seen that in the magnetic unit cell, represented by the dashed line, the chirality sign changes in adjacent plaquettes.

### B. Ferromagnetic Phase Diagram \( \alpha > 0.5 - J < 0 \)

For \( \alpha > 0.5 \), the exchange coupling becomes ferromagnetic \( J < 0 \), \( D \) decreases with \( \alpha \) and the different models are dominated by the intensity of bond anisotropy \( J_{||} \). Let us begin with a very short discussion of the pure F+DM model \( (J_{||} = 0) \). At at low temperatures a typical
phase diagram would show a helical phase at low fields, and a skyrmion lattice emerges at higher magnetic fields. In this situation, two types of intermediate phases may arise: one between the helical phase and the skyrmion crystal, which includes “bimerons”, which look like elongated skyrmions, and another one between the skyrmion crystal and the ferromagnetic phase, a “skyrmion gas”, where the skyrmions are not arranged periodically.

In this work, we explore how the competition of other interactions, fundamentally the presence of the bond anisotropy \( J_\parallel \), combined with the kagome geometry, induces a series of non trivial phases as a function of the external field. We present the low temperature phase diagram for \( \alpha > 0 \) at \( T = 4 \times 10^{-3} \) in Fig. 9. It can be seen, as expected by the presence of strong \( J < 0 \), the system is completely polarized with the external magnetic field for \( \alpha > 0.73 \). However, for lower \( \alpha \), a sequence of non-trivial phases emerge due to the competition between \( J < 0, D \) and \( J_\parallel \).

Starting at low \( \alpha (\alpha \lesssim 0.6) \), the bond anisotropy and the DM interaction are larger than the ferromagnetic exchange coupling (Fig. 2). Taking as an example \( \alpha = 0.55 \), typical textures of the magnetization process are shown in Fig. 10. At low fields (panel A), there is a non trivial “stripe” phase, characterized by a structure factor with two pairs of intense peaks slightly shifted from two \( K \) points in the BZ, and secondary peaks near the third \( K \) point. As the field is increased, there is an intermediate phase where smaller structures start to arise (panel B), and at higher fields a texture reminiscent of a skyrmion crystal is stabilized (panel C). There is a periodic arrangement of small-skyrmion like objects. The structure factor of this phase shows that it is not a skyrmion crystal, since it is a triple-\( q \) phase where \( \sum q \neq 0 \), so we call this phase pSkx-II. We inspect this phase further in panel D, where the spherical snapshot shows a projection of all the spins on the sphere (top view in the inset). Each color represents the tip of the spins from each sublattice. For a skyrmion crystal, the projection covers the whole sphere. Here, two sublattices complete two helices, while a third one forms loops similar to a “Vivani” phase.

For \( \alpha \gtrsim 0.6 \), a remarkable parallel may be drawn between this magnetization process and the one in the F+DM case. Representative snapshots and their corresponding structure factors are displayed in Fig. 11. The lower field phase (panel A) is closer to a helical phase, although the competing interactions and the underlying geometry induce “kinks”. There is a first intermediate phase with “broken spirals”, analog to the bimerons (panel B). Then, the same crystal skyrmion-like structure commented above emerges, with an arrangement of skyrmion-like structures (panel C). Before the system is completely polarized, there is a phase comparable to a skyrmion gas, where these structures appear but do not form a periodic lattice, resulting in a less defined structure factor (panel D).

To complement our study, we resort to the scalar chirality. Fig. 12 shows \( \chi_{nn1} \) (top panel) and \( M_z \) (bottom panel) as a function of \( h_z \) for different \( \alpha \), at \( T = 4 \times 10^{-3} \). It can be seen that, due to their intricated nature, the striped and helical structures have non-zero chirality. The skyrmion-like crystal emerges in the region where \( \chi_{nn1} \) reaches its maximum value and remains almost constant. Finally, the chirality drops at larger magnetic fields (corresponding to the “skyrmion gas”) until the magnetic moments are completely polarized in the saturation state.
FIG. 11: Representative snapshots of typical spin textures for $\alpha = 0.65$ at $T = 4 \times 10^{-3}$, for $h_z = 0.2, 0.5, 0.8; 0.9$ at panels A,B,C,D, respectively. Insets show the corresponding structure factors.

FIG. 12: Scalar chirality (top panel) and magnetization (bottom panel) as a function of the external field at $T = 4 \times 10^{-3}$, for four different values of $\alpha > 0.5$ (ferromagnetic $J$, $\alpha = 0.58, 0.60, 0.65$)

FIG. 13: Panel A: Spin configuration snapshot at $\alpha = 0.5$ and $T = 4 \times 10^{-3}$, $h_z = 3.5$. The inset shows the two types of plaquette order. Panels B,C,D illustrate the three sublattices. Blue arrows in panels B and D indicate the lines of spins that have been “swapped” between these sublattices. The inset in panel C shows the nearest neighbor scalar chirality as a function of temperature for three independent MC realizations at $h_z = 3.5$

FIG. 14: Behavior with temperature for the particular point $\alpha = 0.5$ ($J = 0, J_4 = D$). Panel A: specific heat $C_v$ as function of $T$, for three values of the external magnetic field $h_z = 0.2, 0.45, 0.6$. Three values of $T$ are marked with vertical lines. Panels B-C: Structure factors $s_q^+\perp$ for $h_z = 0.45$ for the three $T$ indicated in panel A.
C. Special case $\alpha = 0.5 - J = 0$

In the particular case where the hopping parameter matches the strength of the Rashba interaction ($t = \lambda$), $\alpha = 0.5$, the exchange coupling is $J = 0$ and $D = J_1 = 0.5$. At low temperatures, we find an “umbrella-like” 120° spin structure with semi-extensive degeneracy, for all the range of external magnetic field. A typical snapshot is shown in Fig. [13] A), with the two types of plaquettes illustrated in the inset. The inspection of the sublattice snapshots reveals a particular form of semiextensive generacy, similar to that seen in the extended Heisenberg model in the kagome lattice[34,35] two sublattices can exchange “lines” of spins, and this gives rise to the two types of umbrella plaquettes, which have opposite chirality. In panels B and D from Fig. [13] blue arrows indicate the lines that have been exchanged between sublattices. Due to its degeneracy, thermal fluctuations select this type of state, and the number and lines swapped between sublattices will change in each independent realization. This will be reflected in the total chirality, as we show in the inset of Fig. [13] panel C, for three MC realizations with $h_z = 3.5$.

In the pure kagome antiferromagnet ($J = 1, J_1 = D = 0$), thermal fluctuations select the most collinear phases at higher field, which produces a pseudoplateau at $M_z = 1/3[23,24]$. To induce the pure “umbrella” configuration in the kagome lattice, interactions such as DM, or further neighbors exchange couplings and $XXZ$ anisotropy are needed[25,26]. Models where this phase is stabilized at low temperature are particularly relevant since they provide a simple magnetic order where interesting transport phenomena may arise[20,21]. In this “umbrella” configuration with semiextensive degeneracy, thermal Hall effect may also be expected, and the transverse thermal conductivity may be a way to explore the state selection.

Moreover, an interesting behaviour with temperature can be seen in the specific heat, $C_v$. In Fig. [14] we plot $C_v$ vs $T$ for three values of the external magnetic field, $h_z = 0.2, 0.45, 0.6$. Coming from higher temperature, there is a first peak. As $T$ is lowered, and the field is increased, a “bump” in the $C_v$ curve can be seen. Inspecting the structure factor, we see that this bump does not seem to be associated to any change in magnetic order. In Fig. [14] panels B-D, we show $S_q^z$ for three values of $T$ in the different regions of the $C_v$ curve for $h_z = 4.5$: before the first peak, between the peak and the bump, and below the bump. For higher $T$, the $S_q^z$ resembles that of the pure kagome antiferromagnet in the cooperative paramagnet phase[25,26], indicating an algebraic spin-liquid behaviour, which is particularly surprising in the absence of isotropic exchange interactions. As $T$ is lowered, there are well defined Bragg peaks at the $M$ points of the EBZ, showing no significant changes for temperatures lower that the bump in the $C_v$.

IV. CONCLUSIONS

The potential applications of antiferromagnetic skyrmions has increased the search for models and materials where these topological structures could be observed. Competing interactions and geometrical frustration seem to be key ingredients in this ongoing exploration. In this work, we have combined these elements, taking an effective spin model with exchange, anisotropic and antisymmetric couplings, on the kagome lattice, which is well known to be highly frustrated in the purely antiferromagnetic case.

First, coming from a microscopic electronic model with Rashba spin orbit coupling where electrons are in contact with a magnetic background via Hund’s couplings, we have derived an effective spin model at half filling for the strong Hund limit, where the parameter $\alpha$ controls the magnitude of the effective exchange, anisotropic and antisymmetric interactions. Then, exploring the low temperature phases with Monte Carlo simulations, we find magnetic configurations fundamentally characterized by their different chiral behavior.

In the antiferromagnetic region of the model ($\alpha > 0.5$), at lower $\alpha$, we find a phase where small locally chiral clusters can be identified. The remarkable difference with the kagome AF+DM model (with $J_1 = 0$) is the absence of the intermediate spiral phase, which implies that the magnetization process is distorted by the $J_1$ interaction, and as the external field increases, a Néel type of pseudo-skyrmion crystal phase is stabilized. The phases that complete the antiferromagnetic $\alpha - h_z$ diagram are nontrivial double-q an single-q phases. The double-q phase is characterized by a net chirality. As the parameter $\alpha$ increases, the anisotropy coupling $J_1$ becomes more relevant, and the single-q phase emerges, where we observe that this phase has no net chirality due an alternation in the sign of the plaquette chirality similar to the so-called “cuboc” phases.

For $\alpha > 0.5$, the exchange coupling is ferromagnetic. Through the calculation of the structure factor and the analysis of the magnetic textures, we find a magnetization process analog to the one that occurs in a pure F+DM model. At low fields, stripe-like and spiral-like phases emerge. Then, at intermediate fields a phase with skyrmion-like textures is stabilized. Moreover, intermediate phases similar to bimeron and skyrmion-gas phases can also be seen.

In the last part of this work, we study the specific point in parameter space ($\alpha = 0.5$) where the exchange interaction $J$ is canceled, while the DM interaction and the anisotropic exchange coupling $J_1$ are equal. In this particular case, we find that thermal fluctuations select an “umbrella” phase with semieextensive degeneracy, where two sublattices can exchange rows of spins without energy cost.

In summary, we have studied an effective model where the combination of geometrical frustration, and competing exchange, DM and anisotropic interactions gives rise
to a rich collection of non-trivial chiral and topological phases. We expect this work to be relevant in the ongoing research of topological phases in frustrated lattices, where much progress is being made in kagome materials.

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