Theory of vortex beam induced skyrmion duplex in chiral magnets

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Application of laser to solids and controlling their electric and magnetic properties are being actively explored in the broad region of condensed matter physics, including spintronics and magneto-optics. We theoretically propose an application of optical and electron vortex beams carrying intrinsic orbital angular momentum to chiral ferro- and antiferro-magnets. We analyze the time evolution of spins in chiral magnets under irradiation of vortex beams, by using Landau-Lifshitz-Gilbert equation. We show that beam-driven nonuniform temperature can induce a new class of topological magnetic defects as well as conventional skyrmions. The new defects have “dounut-shaped” spatial structure of spins which reflects the spatial profile of vortex beams, and therefore cannot be created by ordinary Gaussian beams. We name this new defect as skyrmion duplex and discuss the proper beam parameters and the way of applying the beams for the creation of that. We also study dynamical properties of them under spin-polarized currents and the interaction with magnetic impurities. Our new findings provide an ultrafast way of generating and controlling microscopic magnetic textures in solids with vortex beams and add new ingredients to spintronics and skyrmionics.

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I. INTRODUCTION

Optical vortex, which is electromagnetic wave carrying intrinsic orbital angular momentum (OAM), has been intensively studied in optics since its first proposal in 1992 [1, 2]. More recently, beams of electrons with intrinsic OAM has been also proposed and experimentally realized [3-5]. The OAM of vortex beams is different from the spin degrees of freedom of photons which correspond to the polarization of beams. That is, vortex beams with non-vanishing OAM are different from circularly polarized lights. The intrinsic OAM twists the phase structure of the propagating vortex beams and forces the beams to have topological singularity along the propagation axis, a line with vanishing beam intensity (see Sec. II).

Applications of optical vortices are actively explored. For instance, we can transfer their OAM to classical particles [6-8] or excitons [9] to induce rotational motion of them. Moreover, we can utilize its phase structure for realizing super-resolution microscope [10] or for laser ablation [11-16]. Applications of electron vortex beams have been not so much discussed yet but there is an experimental demonstration of imaging Landau levels of electrons with electron vortex beams [17].

Transfer of OAM or laser ablation by vortex beams can be regarded as “printing” of the spatial structure of vortex beams to physical systems. For example, in Ref. [13-15] it is demonstrated that optical vortices with positive (negative) OAM irradiated to a crystal vaporizes its atoms and as a result, creates a spiral-shaped needle with positive (negative) chirality. As a natural extension of this approach, it is quite interesting to play the same game in electronic or magnetic systems. Indeed, interaction between light and magnets is a hot topic in current condensed matter physics [18-37], especially in spintronics [38] and magneto-optics [19-30]. However, interactions between vortex beams and such microscopic degrees of freedom in solid state materials are not explored well. In this paper, we consider a use of vortex beams for controlling spatial spin texture of solids. To this end, we focus on chiral ferro- and antiferro-magnets.

In a class of chiral magnets, due to Dzyaloshinskii-Moriya (DM) interaction [41, 42], there can appear peculiar, spatially localized magnetic defects, so-called magnetic skyrmions [43] (see Sec. III). Skyrmions in chiral ferromagnets gather growing attentions as a promising candidate for the basis of magnetic memory devices with ultra-low energy consumption [44-47], and very recently, skyrmions in antiferromagnets are also discussed [48, 49]. At present there are many proposals on the creation of skyrmions [50-57] in chiral magnets, some of which are experimentally realized.

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All these proposals basically utilize spatially localized perturbations to “print” skyrmions in the system. That is, chiral magnets can store the information of the external perturbation as topological defects. Therefore, by locally perturbing chiral magnets with vortex beams, intuitively we can create new topological defects reflecting the spatial characteristics of these beams.

On the basis of Landau-Lifshitz-Gilbert (LLG) equation, we confirm that application of vortex beams to chiral ferro- and antiferro- magnets indeed yields such new topological defects with ring-shaped spatial profile, which reflects the spatial profile of the vortex beams. In the case of chiral ferromagnets, the new defect is a bound state of a skyrmion and anti-skyrmion, which we call as skyrmion duplex. In the case of chiral antiferromagnets, there can be a variety of new topological defects other than the simplest, ring-shaped skyrmion duplex. We study basic properties of skyrmion duplexes, their interaction with magnetic impurities and spin-polarized current for both chiral ferro- and antiferro-magnets.

The rest of this paper is organized as follows. In Sec. II we review that vortex beams carry intrinsic OAM and have topological singularity with vanishing beam intensity at their center. We explain how the effect of vortex beams can be theoretically described in magnetic materials. In Sec. III and IV are devoted to our results. We demonstrate the emergence of a new class of topologically stable magnetic defects in these systems other than ordinary skyrmions. We determine optimized way of applying vortex beams for creating these new defects and study how these defects interact with magnetic impurities and spin-polarized currents. Finally, we summarize our results in Sec. V.

II. VORTEX BEAMS

In Sec. II A, we shortly review derivation of optical vortices, or Laguerre-Gaussian (LG) modes of Maxwell’s equations in a vacuum based on the paraxial approximation. We explain that each eigenstate of the equations possesses intrinsic OAM. In Sec. II B we consider how vortex beams interact with magnetic materials. Because of the mismatch in timescales between vortex beams and the dynamics of spins in chiral magnets, we can assume that the effect of vortex beams is heating which realizes nonuniform temperature proportional to the local beam intensity.

A. Optical vortex, Laguerre-Gaussian solutions of Helmholtz equation

Dynamics of electromagnetic fields is described by Maxwell’s equations. In particular, if we assume electromagnetic waves with fixed frequency $\omega$, their propagation in a vacuum is governed by the following wave equations:

$$\left( \Delta + \frac{\omega^2}{c^2} \right) \vec{E} = 0 \quad (1)$$
$$\left( \Delta + \frac{\omega^2}{c^2} \right) \vec{B} = 0, \quad (2)$$

where $\Delta$ is three-dimensional Laplacian and $c$ is the speed of light in a vacuum. These equations are equivalent to Helmholtz-type differential equations.

Vortex beams are defined by solutions of the wave equation

$$\left( \Delta + k^2 \right) \psi(\vec{r}) = 0 \quad (3)$$

in the cylindrical coordinate $(\rho, \phi, z)$, where $\rho$ is the radial coordinate, $\phi$ the azimuthal angle, and $z$ the coordinate along the cylindrical axis. In this coordinate, Laplacian is written as $\Delta = \frac{\partial^2}{\partial \rho^2} + \frac{1}{\rho} \frac{\partial}{\partial \rho} + \frac{1}{\rho^2} \frac{\partial^2}{\partial \phi^2} + \frac{\partial^2}{\partial z^2}$. We assume that the beam propagates along the cylindrical axis. For example, the Maxwell’s equation (2) for the electric field $\vec{E}(\vec{r})$ reduces to Eq. 3 by setting $\frac{\omega^2}{c^2} = k^2$ and $\vec{E}(\vec{r}) = \vec{e} \psi(\vec{r})$ where $\vec{e}$ is the polarization vector. By taking $\psi(\vec{r}) = u(\vec{r}) e^{ikz}$ we have

$$\left( \Delta_T + 2ik \frac{\partial}{\partial z} + \frac{\partial^2}{\partial z^2} \right) u(\vec{r}) = 0 \quad (4)$$

where $\Delta_T = \frac{\partial^2}{\partial r^2} + \frac{1}{\rho} \frac{\partial}{\partial \rho} + \frac{1}{\rho^2} \frac{\partial^2}{\partial \phi^2}$ is the transverse component of Laplacian. Let us employ paraxial approximation to Eq. 4. Namely, we consider the case where $z$-dependence of $u(\vec{r})$ is small in the sense $|\frac{\partial^2 u}{\partial z^2}| \ll |\frac{\partial u}{\partial z}|, |\frac{\partial^2 u}{\partial \rho^2}|$, and $|\frac{\partial^2 u}{\partial \rho^2}| \ll 2k|\frac{\partial u}{\partial \rho}|$ where $k$ is the wavenumber in the $z$-direction. Then we can drop the second order derivative term.
\[
\frac{\partial^2}{\partial z^2} u \text{ in Eq. (4) and have}
\]
\[
\left( \Delta_T + 2ik \frac{\partial}{\partial z} \right) u(\vec{r}) = 0.
\]  

(5)

Vortex beams, or LG modes \[58\] form a complete set of solutions for Eq. (5) and are given by

\[
u^{LG}(\vec{r}) = \frac{1}{\sqrt{|w(z)|}} \left( \frac{\rho}{w(z)} \right)^{|m|} L^{|m|}_p \left( \frac{2\rho^2}{w(z)^2} \right) \times e^{-\frac{\rho^2}{w(z)^2}} e^{-\frac{i\chi(z)(|m|+2p+1)e^{im\phi} e^{-\frac{\rho^2}{w(z)^2}}}{w(z)}},
\]

(6)

where the integers \(p\) and \(m\) label the mode and \(L^{|m|}_p\) is generalized Laguerre function. We note that by setting \(p = m = 0\), Eq. (6) falls into the usual Gaussian beam. The beam width \(w(z) = w\sqrt{1 + \frac{z}{Z}}\) takes its minimum \(w\), called beam waist, at \(z = 0\). Rayleigh range \(Z\) is the distance from the focal plane along the propagation axis at which the cross-section of the beam becomes twice the minimum value. The phase factor determined by \(\chi(z) = \tan^{-1} \left( \frac{z}{Z} \right)\) is called Gouy phase. Since OAM around the propagation axis is given by the eigenvalue of \(L_z = -i\hbar \frac{\partial}{\partial \phi}\), the phase twist \(e^{im\phi}\) in Eq. (6) yields OAM \(\hbar m\). This factor requires the topological singularity \(u^{LG}(0, \phi, z) = 0\) to maintain the single-valuedness of electromagnetic fields at the origin \(\rho = 0\). Hence the OAM leads to “doughnut-shaped” transverse intensity profile of the propagating vortex beams. We emphasize that the OAM is a property of \(u(\vec{r})\) or \(\psi(\vec{r})\) and nothing to do with the polarization of the electromagnetic waves which corresponds to the spin degrees of freedom of photons.

Compared to optical vortex, vortex beam of electrons is more subtle. Here we just mention that in a particular setup, wave function of electrons also acquires the factor \(e^{im\phi}\) with \(m\) an integer (for details, see Ref. [3]).

B. Effect of vortex beams on chiral magnets

In Fig. 1(a) we show schematics of our setup. We place chiral magnets (the figure is for chiral ferromagnets) at the focal plane \((z = 0)\) of vortex beams with ring-shaped spatial profile:

\[
u^{LG}(\rho, \phi, 0) = \frac{\left( \frac{\rho}{w} \right)^{|m|} e^{-\frac{\rho^2}{w^2}} e^{im\phi} L^{|m|}_p \left( \frac{2\rho^2}{w^2} \right) \sqrt{|w|}}{w^{|m|}}.
\]

(7)

In Fig. 2 for several choices of \(p\) and \(m\), we show the beam intensity of the vortex beams \(|u^{LG}(\rho, \phi, 0)|^2\) at \(z = 0\). For \(m = 0\) the beam (7) reduces to the usual Gaussian beam [Fig. 2(a)] whose intensity peaks at the center \((\rho = 0)\). For \(m \neq 0\) as we show in Fig. 2(b)(c)(d), we have topological singularities at the center and the intensity distribution looks like \((p + 1)\)-fold rings.

FIG. 1. (a) Schematics of our setup. Vortex beam is irradiated to a chiral magnet and induces nonuniform temperature proportional to the local intensity of the beam. Arrows represent spins in chiral magnets [see Eq. (8)] and the colors represent their \(z\)-component (red for +1, blue for −1, and green for 0). (b) New topological defect created by the vortex beam (what we call skyrmion duplex). (c),(d) Spin texture of skyrmions in ferromagnetic backgrounds with up and down spins. (e) Spin texture of an anti-skyrmion.
FIG. 2. Intensity of Laguerre-Gaussian modes at the focal plane $|u^{LG}(\rho, \phi, z = 0)|^2$ for several choices of $p$ and $m$. The brightness is proportional to the intensity. For non-vanishing values of $m$, the intensity at the center is exactly zero and there appear $(p + 1)$-fold rings in the intensity distribution. We also see that for larger $m$, the size of rings becomes larger with increasing $m$.

The length scale of magnetic textures in chiral magnets is determined by the exchange coupling and the strength of DM interaction. In particular, the size of skyrmions in them is specified by the ratio of the exchange and DM coupling. If there can exist a new kind of stable defects induced by vortex beams, their size should not be so much different from that of skyrmions. This naive expectation is numerically verified in Sec. III. Then, taking into account the diffraction limit of lasers, the wavelength of vortex beams suitable for our purpose is also close to the size of skyrmions in the target material. Typical size of skyrmions in chiral ferromagnets are $O(1)$ to $O(10)$ nano meters, so that the proper wavelength is in the region of extreme ultraviolet (EUV) lights. Recently, larger skyrmions whose diameter is several hundreds of nano meters or a few micro meters are observed in chiral ferromagnetic films [59–64]. The natural wavelength is then lifted to that of visible lights. In this paper we assume parameters of typical chiral ferromagnets. We note that optical vortex with EUV wavelength can be generated by using nonlinear optical effect or holography [65, 66]. Compared to optical vortex, it is much easier to focus electron vortex beam and achieve this wavelength.

If the wavelength is 10 nm, the frequency of the optical vortex is about 30 Peta Hz. This is much faster than the dynamics of spins in magnetic materials, whose timescale is typically $\sim 1$ Tera Hz. Therefore, the optical vortex with this frequency cannot induce any coherent dynamics of spins. Instead, vortex beams create hot electrons and excite high-energy lattice vibrations which (locally) equilibrate within $100$ fs $\sim 1$ ps [21]. In this work, therefore, we assume that vortex beams only raise the local temperature of the system in accord with the beam intensity. We note that in reality, the temperature gradient may not be so clear. If we use materials with sufficiently large skyrmions, there would be no problems, but if the length scale is in the EUV region (wavelength $\sim 10$ nm), our results would be modified to some extent. Nevertheless, we expect that for our purpose, printing the spatial profile of vortex beams to chiral magnets, details of the temperature distribution is unimportant, as long as the realized temperature reflects the ring-shaped spatial profile of vortex beams. Indeed, we numerically verified that vortex beams simply induce temperature proportional to the local beam intensity : $T(\vec{r}) \propto |u^{LG}(\rho, \phi, 0)|^2$.

The situation for electron vortex beams is more subtle because electrons are charged and the beams generate electromagnetic fields which can affect spin structures of chiral magnets. However, the intensity of electron vortex beams generated by (S)TEM equipments is generically small at present so that the dominant effect will be heating just as optical vortices. Throughout the calculations in this paper, we focus on single-ring vortex beams with $p = 0$ and $m = 5$.

Here we shortly comment on the use of Tera Hz optical vortices whose timescales is comparable with that of spin systems. Using Tera Hz optical vortices, we can control dynamics of spins by electromagnetic fields themselves, not by heating caused by them. Due to their large length scale (300$\mu$m for 1 Tera Hz), the topological singularity of Tera Hz optical vortices would not play essential roles in that case, but the twisted phase structure of the beams may offer a new way to control large-scale, magnetic domain structures in solids. Moreover, recently it is becoming possible to create vortex beams with their waist $w$ much smaller than their wavelength by using plasmonics techniques [67, 68]. With such techniques, it may be possible to study non-thermal effect of Tera Hz optical vortices reflecting the spatial structure of the beams [69].
III. CHIRAL FERROMAGNET

In this section, we discuss topological defects induced by vortex beams in chiral ferromagnets (FMs) and their basic properties. First we shortly review static properties of chiral FMs, the phase diagram and characteristic magnetic defects, skyrmions in a canonical model of them. Then, we explain our numerical calculations based on (stochastic-) Landau-Lifshitz-Gilbert equation. We show the emergence of new topological defects created by vortex beams and optimize the way of applying the beams. Finally, we discuss how the new defects interact with magnetic impurities and spin-polarized currents.

![Phase Diagram](image)

FIG. 3. Ground state phase diagram of the canonical model of chiral ferromagnets (8) (reproduced from Ref. 70). There appear three distinct phases depending on the magnitude of the external magnetic field applied in the $z$-direction. For the helical order phase and skyrmion crystal phase, we give their spin texture obtained by numerical simulations based on LLG equation, using arrows and colors.

A. Static properties of chiral magnets

Chiral FMs such as B20-type compounds MnSi \(71,74\), Fe\(_{1-x}\)Co\(_x\)Si \(75,76\), or FeGe \(62,77,78\), are characterized by their relatively large DM interaction. A canonical model of chiral FMs is the following two-dimensional classical spin model defined on a square lattice with $\vec{m}_\vec{r}$ being the spin at the site $\vec{r}$:\(8\):

$$H = -J \sum_{\vec{r}} \vec{m}_\vec{r} \cdot \left( \vec{m}_{\vec{r}+a\vec{e}_i} + \vec{m}_{\vec{r}+a\vec{e}_j} \right) + \sum_{\vec{r}} \vec{D}_i \cdot \left( \vec{m}_\vec{r} \times \vec{m}_{\vec{r}+a\vec{e}_i} \right) - B_z \sum_{\vec{r}} m^z_\vec{r},$$

(8)

where $a$ is the lattice constant and $\vec{e}_i$ is the unit vector along the $i$-axis ($i = x, y$). The exchange coupling $J > 0$ represents ferromagnetic Heisenberg interaction and $\vec{D}_i$ is DM vector on the bond $(\vec{r}, \vec{r} + a\vec{e}_i)$, and $B_z$ is the external magnetic field applied in the $z$-direction. Hereafter we normalize the length of $\vec{m}_\vec{r}$ to unity. As long as we are interested in physics like spin waves or magnetic defects whose length scale is much longer than the lattice constant, we expect that the model (8) works as a standard model of chiral ferromagnets. Indeed, many experimental results on real materials are well described by this model (see for example Ref. 80), even though their microscopic Hamiltonians must be much more intricate.

What is peculiar about chiral magnets is the emergence of topologically stable magnetic defects, skyrmions. In Fig. (d), we show skyrmions for a particular choice of DM vectors $\vec{D}_i = D\vec{e}_i$ ($D > 0$) and the external magnetic field $B = \pm B_z \vec{e}_z$. In Fig. (e) we also show the spin texture of a so-called anti-skyrmion. (Anti-) skyrmions are characterized by their non-vanishing skyrmion number:

$$N_{SK} = \frac{1}{4\pi} \int \vec{m}_\vec{r} \cdot \left( \frac{\partial \vec{m}_\vec{r}}{\partial x} \times \frac{\partial \vec{m}_\vec{r}}{\partial y} \right) d^2r.$$  

(9)

In a continuous space, $N_{SK}$ must be quantized to an integer, so that (anti-) skyrmions are stable against any continuous deformations. Even for lattice systems, where the quantization is incomplete, skyrmions and anti-skyrmions are
topologically stable if their size is much larger than the lattice constant. Nevertheless, in chiral FMs, anti-skyrmions are energetically unstable \cite{81} and have short lifetime for the given DM vectors $\vec{D}_i = D\vec{e}_i$.

The phase diagram of the model \cite{8} with $\vec{D}_i = D\vec{e}_i$ is well studied (see for example Ref. \cite{70}). For a fixed exchange and DM interaction, when the external magnetic field is very small, the ground state develops a helical magnetic order with a single wavelength determined by the magnitude of the DM interaction. As we increase the magnetic field, the helical order becomes unstable and the ground state turns into a triangular lattice of skyrmions. For very large magnetic field, spins become polarized in the direction of the field and the ferromagnetic ground state is stabilized. Therefore, the phase diagram of this model is summarized as Fig. 3. We present the critical values of the magnetic field and spin textures of typical states in the helical order phase and skyrmion crystal phase. In addition to the skyrmion crystal phase, skyrmions can appear as energetically stable isolated magnetic defects in the ferromagnetic phase.

![Image of skyrmion duplex](image)

**FIG. 4.** Skyrmion duplex obtained by numerical simulation of the model \cite{8} with $J = 1$, $D = 0.15$, and $B_z = 0.015$. (a) Cumulative skyrmion number $N_{\text{SK}}(R)$ as a function of $R/a$, where $a$ is the lattice constant. The skyrmion duplex is a skyrmion with positive $N_{\text{SK}}$ [Fig. 1(d)] surrounded by that with negative $N_{\text{SK}}$ [Fig. 1(c)]. (b) Spin texture of the skyrmion duplex. The arrows represent the in-plane components of spins and the color indicates their z-component.

### B. Skyrmion duplex

With numerical calculations based on stochastic LLG (sLLG) equation, we confirm that the ring-shaped temperature profile induced by vortex beams offers a way to create a new topological magnetic defect shown in Fig. 1(b) and Fig. 1(b). Details of numerical methods and optimal way of applying vortex beams will be discussed in Sec. III C.

To see the topological nature of this defect, in Fig. 1(a), we present the cumulative skyrmion number

$$N_{\text{SK}}(R) = \frac{1}{4\pi} \int_{r < R} \vec{m}_r \cdot \left( \frac{\partial \vec{m}_r}{\partial x} \times \frac{\partial \vec{m}_r}{\partial y} \right) d^2 r.$$  \hspace{1cm} (10)

The integration is performed within the circle with radius $R$ measured from the center of the defect. The spin texture and the $R$ dependence of $N_{\text{SK}}(R)$ clearly show that the new, ring-shaped defect is a bound state of two skyrmions in Fig. 1(c)(d) where the former (latter) has $N_{\text{SK}} = -1 (+1)$ if placed in a continuous space and isolated. Namely, its spin configuration outside a certain radius, at which $m^z = -1$, is that of the skyrmion in Fig. 1(c) and the spins inside that radius form the skyrmion in Fig. 1(d). Reflecting this structure, we name this new defect as skyrmion duplex (SkD). SkDs have vanishing skyrmion number in total. We find that an SkD is a metastable excited state in the skyrmion crystal phase of the present model, which has longer lifetime than the periods of our LLG calculations.

Spin structures in magnets like soliton lattices and skyrmion lattices can be directly observed by Lorentz TEM \cite{82,83}. Therefore, if SkDs in real materials have sufficiently long lifetime as predicted by our LLG calculations, they should be also experimentally observable using Lorentz TEM.
C. Optimal way of applying vortex beams

Here we determine the optimal way of applying vortex beams to achieve high probability of creating SkDs. Our numerical calculations are based on sLLG equation \[84\] for the model \[8\]:

\[
\frac{d\vec{M}_r}{dt} = -\gamma\vec{M}_r \times \left( -\frac{\partial H}{\partial \vec{M}_r} + \vec{h}_{T(r)}(t) \right) + \alpha \vec{M}_r \times \frac{d\vec{M}_r}{dt},
\]

where \( \vec{M} = h / \gamma \vec{m} \) and \( \gamma \) is gyromagnetic ratio. The time coordinate is \( t \). The second term in the right hand side in Eq. \[11\] is so-called Gilbert damping term which describes dissipation. The dimensionless constant \( \alpha \) characterizes the strength of the dissipation. In the framework of the sLLG equation, the effect of heating is treated as a random field \( \vec{h}_{T}(t) \) satisfying

\[
\langle h_{T(r)}^{\mu}(t) \rangle = 0,
\]

\[
\langle h_{T(r)}^{\mu}(t) h_{T(r')}^{\nu}(t') \rangle = \sigma(\vec{r})\delta^{\mu,\nu}\delta(\vec{r} - \vec{r}')\delta(t - t'),
\]

where \( \mu, \nu = x, y, z \). Here \( \sigma(\vec{r}) \) is determined by temperature at \( \vec{r} \) from the fluctuation-dissipation theorem: \( \sigma(\vec{r}) = 2k_B T(\vec{r}) \alpha \). Therefore, by using Eq. \[11\] with this random field we can numerically simulate the time evolution of the model \[8\] at finite temperatures.

With the fixed parameters \( J = 1, D = 0.15, \) and \( \alpha = 0.1 \), we look for the optimal way of creating SkDs by changing other parameters. We focus on the skyrmion crystal phase but the initial state of simulations \( t = 0 \) is taken as a metastable, perfect ferromagnetic state i.e. \( \vec{m}_r = (0, 0, 1) \) for all \( \vec{r} \). For \( J \) and \( D \), the phase boundary between the helical order phase and the skyrmion crystal phase is \( B_z = 0.0052 \) and that between the skyrmion crystal and the ferromagnetic phase is \( B_z = 0.018 \). After several trials and errors we find that “annealing” the spins is advisable to obtain SkDs with high probability. Hence we fix the time dependence of the vortex-beam driven temperature as shown in Fig. 7(d): \( T(t) = T_0 \left( 1 - \frac{t - t_0}{T_0} \right) \Theta(t)\Theta(t_0 - t) \) with \( t_0 = 500 \). Here \( \Theta(x) \) is the step function. Namely, we assume that the temperature is instantaneously raised to its maximum \( (T_0) \) and gradually cooled down. The time is measured in the unit of \( h/J \), which is estimated to be 0.7 ps for \( J = 1 \) meV. For this strength of \( J \), \( B_z = 0.01 \) corresponds to 0.173 Tesla and \( T = 1 \) does 11.6K. We renormalize the beam intensity by its peak value so that the temperature profile is given by \( T(t, \vec{r}) = T(t) \left( |u_{LG}(\rho, \phi, 0)|^2 / \max_{\vec{r}}(|u_{LG}(\rho, \phi, 0)|^2) \right) \). As we noted in Sec. II B, we take the single-ring vortex beam with \( p = 0 \) and \( m = 5 \). In Fig. 5 we present a particular trial which eventually ends in an SkD. We see that the spin texture perturbed by the vortex beam is annealed to form an SkD. We try the simulations 20 times for each set of \( (T_0, B_z, w) \) in a periodic system with \( 150 \times 150 \) sites and calculate the success probability of creating SkDs. We take the peak temperatures \( T_0 = 1, 1.5, \) and 2. We use Heum method with the time step \( \Delta t = 0.02 \) for numerical integration of the sLLG equation with linearization technique \[80\].

![FIG. 5. Time evolution of the z-component of spins for a particular trial with parameters \( B_z = 0.1, D = 0.15, 150a/w = 12, \) and \( T_0 = 2 \). We start from the metastable ferromagnetic state at \( t = 0 \). At \( t = 0 \) we suddenly raise the temperature in accord with the intensity of the vortex beam with \( p = 0 \) and \( m = 5 \). The the temperature is lowered gradually and finally, a skyrmion duplex is formed. The time is measured in the unit of \( 1/(J\gamma) \). When \( J = 1 \) meV the time unit corresponds to 0.7 ps.](attachment:image.png)

In Fig. 7 we show the success probability of creating SkDs, \( (N/20 \times 100) \% \), where \( N \) is the number of success trials in which we obtain SkDs, and 20 is the total number of trials for each set of \( (T_0, B_z, w) \). The probability tops when the magnetic field is small and the beam waist satisfies \( 150a \sim 13w \). If we take \( w = 0.5 \lambda \) and \( a = 5 \lambda \), as the
FIG. 6. Time evolution of the $z$-component of spins for a particular trial with parameters $B_z = 0.1$, $D = 0.15$, and $T_0 = 1.5$. We still use vortex beam with $p = 0$ and $m = 5$ but the beam waist is set to be small $150a/w = 20$. Due to the small beam waist, the topological singularity in the temperature profile is smeared and as a result, an ordinary skyrmion is created.

We numerically confirm that the size of SkDs grows up with decreasing $B_z$ and it is impossible to create stable SkDs in the ferromagnetic phase. As we decrease the magnetic field $B_z$, the ferromagnetic initial state we assume would become more unstable in reality, while SkDs themselves are easier to create and energetically stabler.
FIG. 8. Time evolution of a skyrmion duplex under the spin-polarized current \( v_x = -0.01 \) with and without a magnetic impurity (emphasized as a white dot), assuming \( \beta = \alpha \).

(a) Initial state at \( t = 0 \) with the magnetic impurity. (b) Initial state without the impurity. The skyrmion duplex is pinned by the impurity. (d) Final state without the impurity. We place white lines indicating the initial position of the skyrmion duplex for guide to eyes.

The calculations are performed for a periodic system with size 300 \( \times \) 300 sites, and the spin texture of a subregion with size 200 \( \times \) 150 sites is shown as colored pixels following the color bar.

FIG. 9. Collision between a skyrmion duplex and magnetic impurity driven by spin-transfer torque under the condition \( \beta = \alpha \).

(a) Snapshots for \( v_x = 0.01 \) at \( t = 0, 2000, 6000, \) and 9600. (b) Snapshots for \( v_x = 0.05 \) at \( t = 0, 200, 600, \) and 800.

D. Dynamical properties of skyrmion duplex

We demonstrated that by applying vortex beams to chiral FMIs, we can create SkDs. Here we study dynamical properties of SkDs in metallic chiral FMIs where localized spins \( \vec{m}_r \) couple to conducting electrons through the s-d interaction. We numerically investigate how they interact with magnetic impurities and in-plane spin-polarized currents flowing in the \( x \)-direction. In the presence of spin-polarized currents, spins feel spin transfer torque (STT) \[85\] with which we can control the center-of-mass motion of skyrmions \[86\]. LLG equation including STT at \( T = 0 \) is

\[
\frac{d\vec{M}_f}{dt} = -\gamma \vec{M}_f \times \left( -\frac{\partial H}{\partial \vec{M}_f} \right) + \frac{\alpha}{M} \vec{M}_f \times \frac{d\vec{M}_f}{dt} + \frac{qa^2}{2|e|} \left( \vec{r} \cdot a\vec{\nabla} \right) \vec{M}_f - \frac{qa^2}{2|e|} \beta \vec{M}_f \times \left( \vec{r} \cdot a\vec{\nabla} \right) \vec{M}_f.
\] (13)

Here \( \vec{r} = j^p \hat{e}_z \) is the density of the electric current with spin-polarization \( q \) (for simplicity we set \( q = 1 \) in the following) along the \( z \) axis, and \( e \) is the elementary charge. The final term with the dimensionless parameter \( \beta \) represents non-adiabatic effect of spin-polarized currents. We define \( j_{0}^{FM} = \frac{2|e|q}{a^2} \) and introduce \( v_x = j^p / j_{0}^{FM} \) to represent the amount of the current. For \( J = 1 \) meV and \( a = 6 \AA \), we have \( j_{0}^{FM} = 1.35 \times 10^{12} \) A/m\(^2\). We assume that temperature is negligibly small and solve Eq. (13) using fourth-order Runge-Kutta method with numerical time step \( \Delta t = 0.2 \) in a periodic system with 300 \( \times \) 300 sites. The parameters are set to be \( J = 1, D = 0.15, \alpha = 0.1, \) and \( B_z = 0.015 \).
The motion of (anti-)skyrmions under spin-polarized currents can be decomposed into longitudinal and transverse motions \([N_0]\). The direction of the transverse drift depends on the sign of the skyrmion number. Since an SkD is a composite of two topological defects with \(N_{SK} = \pm 1\), the transverse drift is expected to cause deformation of SkDs. When \(\beta = \alpha\), the transverse motion is known to cancel out \([N_0]\). Hence, for the study of the center-of-mass motion of SkDs in the presence of magnetic impurities, it is convenient to assume \(\beta = \alpha\). In realistic materials this is not always a reasonable assumption, and the deformation would be experimentally relevant unless the amount of the current is very small. The motion of skyrmions can be also driven by other mechanisms such as thermal gradient \([Z7]\). If there is no transverse motion depending on \(N_{SK}\) driven by some mechanism, our results in the following will qualitatively apply to that situation. The outcome of the deformation of SkDs will be discussed later.

We consider a single site magnetic impurity with fixed spin direction \(\vec{s} = (0, 0, 1)\) and examine the interplay between an SkD and the impurity. For simplicity, the interaction Hamiltonian between the impurity and other spins is set to be the same as that among other spins. First, we set a magnetic impurity at a midpoint of an SkD. Since an SkD is a kind of ring-shaped domain wall, we expect that it will be pinned by such magnetic impurity under spin-polarized currents. To confirm this expectation, we compare the time evolution of an SkD in the presence and absence of the magnetic impurity as shown in Fig. 8 where the impurity is emphasized as a white dot. We apply electric current \(v_x = 0.01\) and follow the time evolution of the SkD. For visibility we place white lines in the figures to indicate the initial position of the SkD. Starting from the initial state at \(t = 0\) [Fig. 8(c)], the SkD is driven by STT until \(t = 4000\). In the presence of the magnetic impurity [Fig. 8(c)], the SkD remains in the initial position and indeed is pinned by the impurity. On the other hand, without the impurity the SkD moves as shown in Fig. 8(b(d)).

Next we consider collision between SkDs and impurities under the condition \(\beta = \alpha\). As Fig. 9(a) shows, when their relative velocity is sufficiently small, the SkD avoids the impurity. On the other hand, if the relative motion is fast, the SkD goes over the defect by partially breaking its ring-shaped structure. After the penetration process, the spin texture is recovered to that of an SkD as is shown in Fig. 9(b). Therefore, SkDs can move around chiral FMs with their spatial structure kept in the presence of magnetic impurities and spin-polarized currents under the condition \(\beta = \alpha\).

Finally we examine the role of the transverse drift which we ignored above by setting \(\beta = \alpha\). As we noted, the transverse motion offers a way to realize continuous deformation of SkDs. Since an SkD has vanishing skyrmion number, there are two possibilities after the deformation. The first possibility is to obtain a trivial ferromagnetic state without any topological defects. The other possibility is to have a spatially separated pair of a skyrmion and anti-skyrmion in the ferromagnetic background. We numerically simulate the time evolution of an SkD under the spin-polarized current \(v_x = -0.1\Theta(3500 - t)\Theta(t)\) in a periodic system with \(300 \times 300\) sites. The parameters are taken to be \(J = 1, D = 0.15, \alpha = 0.1, B_z = 0.015\), and \(\beta = 0\). As shown in Fig. 10(a), an SkD deforms and breaks off into a spatially separated pair of down-spin domains. These two domains are indeed a skyrmion (upper one) and anti-skyrmion (lower one) as is confirmed from the comparison of Fig. 1(c)(d) and Fig. 11. The anti-skyrmion in Fig. 11 gradually shrinks to vanish as we see in Fig. 10(b) since it is energetically unstable. Since the anti-skyrmion has positive skyrmion number \(N_{SK} = +1\), that of the whole system changes from zero to minus one when the anti-skyrmion vanishes. Figure 12 shows that at \(t = 26000\), the skyrmion number decreases by one. Besides, Fig. 10(a) shows that the breakdown of the SkD completes at \(t = 1000 \sim 1500\). Therefore, the lifetime of the anti-skyrmion is estimated to be about \(\tau \sim 25000/(\gamma J)\), which corresponds to 18 ns if \(J = 1\) meV.

In a previous study \([54]\) of local heating by near-field, only skyrmions are observed in chiral ferromagnets. Another theoretical study \([81]\) shows that by magnetic field quenching on helical ordered states, we can create anti-skyrmions. However, the number and position of these anti-skyrmions are uncontrollable, and they have too short lifetime to be experimentally observed by Lorentz TEM. On the other hand, our result offers a unique scheme to create an anti-skyrmion in chiral FMs at desired position with relatively long lifetime. The long lifetime of the anti-skyrmion originates from its large initial size and we expect that we can enhance the lifetime by starting from a larger SkD, by decreasing \(B_z\) or using proper materials like CoFeB-Ta.

**IV. CHIRAL ANTIFERROMAGNET**

Next we move onto chiral antiferromagnets (AFMs). The basic idea is the same as the ferromagnetic case. Namely, we consider of printing the spatial structure of vortex beams to chiral antiferromagnets as topological defects. We model the application of vortex beams as heating which realizes temperature proportional to the local intensity of the beams, and examine the time evolution of spins by sLLG equation.

We first review the phase diagram of a canonical model of chiral AFMs. Then we show the emergence of antiferromagnetic analogue of SkDs after the laser irradiation and discuss how to achieve high success probability of their creation. Finally we study dynamics of antiferromagnetic SkDs in the presence of magnetic impurities and spin-polarized currents.
FIG. 10. Continuous deformation of a skyrmion duplex by spin transfer torque from the uniform spin-polarized current $v_x = -0.1\Theta(t)(3500 - t)$. We show snapshots of the $z$-component of spins $m^z_{\vec{r}}$ at (a) $t = 0, 1000, 1500,$ and $2500$. (b) Snapshots at $t = 5000, 15000, 20000,$ and $26000$.

FIG. 11. Spin structure at $t = 12000$ within the process of Fig. 10. By comparing the spin structure of the upper (lower) defect with that of the (anti-) skyrmion in Fig. 1(c) [Fig. 1(e)], we confirm that the skyrmion duplex is decomposed into a skyrmion (upper one) and anti-skyrmion (lower one).

A. Static properties

We use the following canonical model of chiral antiferromagnets on a square lattice

$$H_{AF} = J \sum_{\vec{r}} \bar{m}_{\vec{r}} \cdot \left( \bar{m}_{\vec{r}+a\vec{e}_x} + \bar{m}_{\vec{r}+a\vec{e}_y} \right) + \sum_{\vec{r}} \vec{D}_i \cdot \left( \bar{m}_{\vec{r}} \times \bar{m}_{\vec{r}+a\vec{e}_i} \right) - A \sum_{\vec{r}} (m^z_{\vec{r}})^2, \quad (14)$$

where a spin localized at $\vec{r}$ is represented as $\bar{m}_{\vec{r}}$ with its norm normalized to unity. The exchange coupling is antiferromagnetic ($J > 0$) and $A$ is the anisotropy energy along the $z$-axis. The DM interaction stabilizes antiferromagnetic skyrmions in this model. For example, if we choose $\vec{D}_x = D\vec{e}_y$, $\vec{D}_y = -D\vec{e}_x$, Néel-type antiferromagnetic skyrmion [shown in Fig. 13(a)] appears. Basic properties of these Néel-type antiferromagnetic skyrmions, such as their dynamics under applied spin-polarized currents and their stability at finite temperature are studied recently [48, 49].

When $A/J = 0.055$, the phase diagram of this model [49] with $\vec{D}_x = D\vec{e}_y$, $\vec{D}_y = -D\vec{e}_x$ is given in Fig. 14. For $D/J < 0.22$, the ground state is a Néel state, but we can divide this phase into two regions. In the region with $D/J < 0.16$, antiferromagnetic skyrmions are unstable while if $D/J > 0.16$, they are energetically stable as isolated topological defects with very long lifetime. Following Ref [49], we call the latter region as antiferromagnetic skyrmion (AFMS) region. For larger DM interaction, skyrmions deform (d-AFMS region) and eventually turn into warm domains (WD region). In lower panels of Fig. 14 we show typical states in these phases as staggered spins $m^z_{\vec{r}} \times (-1)^{|\vec{r}|} \equiv m^z_{\bar{r}=(i,j)} \times (-1)^{i+j}$. We note that WDs can be regarded as a collection of strongly deformed skyrmions, so that the phase boundary between the d-AFMS phase and WD phase is unclear from our LLG calculations in a finite size system.
FIG. 12. Time dependence of the skyrmion number within the process of Fig. 10. We see that the skyrmion number suddenly changes at $t = 26000$ by one. It corresponds to the disappearance of the anti-skyrmion.

FIG. 13. Schematics of a Néel-type antiferromagnetic (a) skyrmion and (b) skyrmion duplex. There are two magnetic sublattices in Néel ordered states in the square lattice, and a skyrmion and skyrmion duplex can be seen as bound states of their ferromagnetic counterparts living in different magnetic sublattices. Just as Fig. 1, the colors of arrows represent the $z$-component of spins.

B. Antiferromagnetic skyrmion duplex

We focus on the AFMS region in Fig. 14 and take $\vec{D}_x = D\vec{e}_y$, $\vec{D}_y = -D\vec{e}_x$ in the following. We treat the effect of applied vortex beams by the sLLG equation, assuming temperature proportional to the beam intensity. Then we observe Néel-type antiferromagnetic SkDs [Fig. 13(b)] after the irradiation in the proper manner (see Sec. IV C). Observing the spin texture carefully, we find that an antiferromagnetic SkD consists of two ferromagnetic SkDs in the two magnetic sublattices. Antiferromagnetic SkDs are energetically stable and their lifetime is longer than our calculation periods. In the case of anti-ferromagnets, experimental observation of their spin textures with Lorentz TEM is difficult. In Ref. 48, instead, neutron scattering and X-ray magnetic linear dichroism are proposed as a probe of antiferromagnetic skyrmions. We expect that these methods are also applicable to antiferromagnetic SkDs.

C. Optimal way of applying vortex beams

Following the analysis in the ferromagnetic case, here we optimize the way we apply vortex beams to achieve high success probability of creating antiferromagnetic SkDs. The sLLG equation for the antiferromagnetic Hamiltonian [14]
FIG. 14. Phase diagram of the canonical model of chiral antiferromagnets Eq. (14) for $A/J = 0.055$. When DM interaction is very weak, we have antiferromagnetic (AFM) region where we cannot have skyrmions. As we increase DM interaction $D$, antiferromagnetic skyrmions become energetically stable at $D/J \sim 0.16$ as isolated defects. For larger DM interaction ($D/J \geq 0.22$), skyrmions deform to lower their energy (d-AFM state) and eventually warm domains are formed (WD). The phase boundary between d-AFM and WD is unclear from our calculations. We visualize the spin texture of typical states in each phase obtained from LLG calculations by using staggered spins $m^z_r \times (-1)^{|\bar{r}|} \equiv m^z_{\bar{r}=m(i,j)} \times (-1)^{i+j}$.

FIG. 15. Time evolution of staggered spins $m^z_r \times (-1)^{|\bar{r}|} \equiv m^z_{\bar{r}=m(i,j)} \times (-1)^{i+j}$ for a particular trial with parameters $D/J = 0.205$, $A/J = 0.055$, $T_0/J = 1$, and $m = 5$ creates magnetic domains different from the background Néel order. These domains merge to form an antiferromagnetic skyrmion duplex. A time step corresponds to 0.7 ps for $J = 1$ meV.

The Hamiltonian is

$$\frac{d\vec{M}_F}{dt} = -\gamma \vec{M}_F \times \left( -\frac{\partial H_{AF}}{\partial \vec{M}_F} + \vec{h}_{T(\bar{r})}(t) \right) + \alpha \vec{M}_F \times \frac{d\vec{M}_F}{dt}, \quad (15)$$

where $\vec{M}_F = h\gamma \vec{m}_F$, and $\vec{h}_{T(\bar{r})}(t)$ is again the random field satisfying Eq. (12) and $\sigma(\bar{r}) = 2k_B T(\bar{r})\alpha$. We take a system with $150 \times 150$ sites and impose periodic boundary condition. The time step of Heun method is set to $\Delta t = 0.03$. Hereafter we take $J = 1$ and fix $D = 0.205$ and $A = 0.055$. After several trials and errors, we find that it is advantageous to keep the temperature constant for long periods to achieve high probability. Therefore we assume the time dependence of the temperature induced by vortex beams as $T(t) = T_0 \Theta(t_0 - t) \Theta(t)$ [shown in Fig. 16(d)] and examine the probability for several values of $t_0 = 3000, 5000,$ and $7000$. Here $T_0$ is set to be proportional to the local intensity of vortex beams with $p = 0$ and $m = 5$, and $\Theta(x)$ is the step function. In Fig. 15 we show time evolution within a particular trial. We see that under the static heating, domains of a Néel state different from the background one grow. These domains merge and result in an antiferromagnetic skyrmion duplex. We try calculations 20 times for each set of $(T_0, t_0, w)$ and obtain the success probability in the same way as the ferromagnetic case. The numerical calculations are performed in the AFMS region of chiral AFMs where the ground state is a Néel ordered state. Therefore, we choose a stable Néel state as the initial state of our calculations. This is in contrast with the ferromagnetic case where we assumed the metastable ferromagnetic initial state. The obtained
The success probability is summarized in Fig. 16. The optimal beam waist is found to satisfy $150a/w \sim 13$. If the lattice constant is $a = 5\text{Å}$ and the beam focused well $\lambda \sim 2w$, the wavelength $\lambda \sim 13\text{nm}$, again close to the size of (antiferromagnetic) skyrmions and in the EUV region. Just as the ferromagnetic case, we expect that the wavelength can be lifted by using materials with larger antiferromagnetic skyrmions.

Contrary to the ferromagnetic case, there appear various kinds of topological defects other than the simplest, ring-shaped antiferromagnetic SkDs. In Fig. 16(e) we show $z$-component of staggered spins for several examples of topological defects we observed, including the simplest SkD. In the case of chiral FMs, due to the external magnetic field $B_z$, complicated magnetic structures are energetically unfavored and only the simplest SkD is observed within our calculations. However, since there is no corresponding “staggered magnetic field” in the present model, two possible Néel states are energetically degenerate. Therefore, even if the spin texture of a topological defect is intricate, it costs energy only at the domain boundaries and can be energetically stable. All these defects are roughly speaking, ring-shaped, and reflect the spatial profile of vortex beams. We confirm that they have lifetime longer than our calculation periods. Therefore, in making Fig. 16 we regard a trial to be in success if one of such ring-shaped defects are observed after the irradiation of vortex beams.

Throughout our calculations we have assumed the simplest vortex beams with $p = 0$ where the spatial profile of their intensity looks like a single ring. However, as Fig. 2 shows, by changing the integer parameter $p$ in Eq. (7), we can realize vortex beams with spatial profile of $(p+1)$-fold rings. Since topological defects in chiral AFMs do not cost large energy as we noted, we expect that changing the parameter $p$ of vortex beams can further increase the family of possible topological defects in chiral AFMs.

**FIG. 16.** Success probability of creating antiferromagnetic skyrmion duplexes by vortex beams with $p = 0$ and $m = 5$. We fix $J = 1$, $D = 0.205$, $A = 0.055$, and $\alpha = 0.1$. The initial state at $t = 0$ is a Néel ordered ground state and the temperature is varied in accord with (d) for (a) $t_0 = 3000$, (b) $t_0 = 5000$, and (c) $t_0 = 7000$. The probability is high when the wavelength of vortex beams is of the same order of the size of antiferromagnetic skyrmions, and the period of the irradiation of vortex beams is long. (e) Several examples of topological defects observed after the irradiation of vortex beams. We show $z$-component of staggered spins for visibility.

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### D. Dynamical properties of antiferromagnetic skyrmion duplex

We study dynamics of antiferromagnetic SkDs under spin-polarized currents and magnetic impurities, focusing on metallic chiral AFMs. The only difference from the ferromagnetic case is the way STT is introduced. Since the spin texture of AFMs is not spatially smooth, we cannot use the previous formulation of STT caused by in-plane currents, which are described by terms with spatial derivative [see Eq. (13)]. Instead, here we assume that electric currents are injected perpendicularly to the $x$-$y$ plane and their polarization is along the $-y$ direction [49, 52]. In this case, the effect of STT within the framework of LLG equation is described by so-called Slonczewski term. The LLG-Slonczewski...
where \( j \) is the density of the spin-polarized current with polarization \( q \) (for simplicity we assume \( q = 1 \) in the following), and the polarization direction \( \vec{p} = -\hat{y} \). The last term in the right hand side is the Slonczewski term. Here we assume that the system is a thin film of thickness \( d \) with saturation magnetization \( M_s \). Assuming that the film is made of two-dimensional layers stacked in the \( z \)-direction and inter-layer interactions are negligible, we apply Eq. (16) to the two-dimensional Hamiltonian Eq. (14). We define \( j_0^{\text{AFM}} = \left( -\frac{2e|\vec{M}|d}{\hbar q} \right) \) and use \( v_x = j/j_0^{\text{AFM}} \) to measure the current. For \( M_s = 290 \text{kA/m} \), \( d = 0.4 \text{nm} \), and \( J = 1 \text{meV} \), we have \( j_0^{\text{AFM}} \approx 3.04 \times 10^{12} \text{ A/m}^2 \), which is of the same order of the corresponding quantity in the ferromagnetic case \( j_0^{\text{FM}} = \frac{2e|\vec{M}|}{\hbar q} = 1.35 \times 10^{12} \text{A/m}^2 \) for \( a = 5\text{Å} \).

The motion of antiferromagnetic skyrmions induced by the Slonczewski term is studied in Ref. [49] and [52]. Since an antiferromagnetic skyrmion can be regarded as a tight bound state of two ferromagnetic skyrmions living in different magnetic sublattices as shown in Fig. [13(a)], there is no analogue of the transverse motion in the case of skyrmions in chiral FMs. As a result, the motion of antiferromagnetic skyrmions and SkDs under the spin-polarized currents are the same. Actually, we can numerically verify that for \( \vec{p} = -\hat{y} \), both antiferromagnetic skyrmions and SkDs move in the \( x \)-direction.

Using Eq. (16), we numerically calculate dynamics of antiferromagnetic SkDs in the presence of magnetic impurities and spin-polarized currents in a periodic system with \(200 \times 200\) sites. We assume that the magnetic impurities are pinned in the direction consistent with the background Néel state. As shown in Fig. [17] and [18], antiferromagnetic SkDs exhibit qualitatively the same behaviors against magnetic impurities and STT as ferromagnetic SkDs do. Namely, antiferromagnetic SkDs are pinned by impurities (Fig. [17]) and scattered by impurities in different ways, depending on the speed of antiferromagnetic SkDs (Fig. [18]). We also see that antiferromagnetic SkDs are easily deformed by impurities and currents. This is attributed to the fact that there is no “staggered magnetic field” stabilizing one of Néel states.

The magnitudes of spin-polarized currents used in the calculations above are much smaller than those in the ferromagnetic case, even though we have qualitatively the same results. The previous study in Ref. [49] also observes the same feature for ordinary skyrmions. Namely, antiferromagnetic skyrmions can be driven with much smaller amount of current than ferromagnetic skyrmions. Since the center-of-mass motion of (antiferromagnetic) SkDs is the same as the longitudinal motion of isolated (antiferromagnetic) skyrmions, the quantitative difference between ferromagnetic and antiferromagnetic SkDs we observed can be understood from that of skyrmions.

In summary, we see that we can create antiferromagnetic SkDs by vortex beams and control their motion with small amount of electric currents. We confirm that antiferromagnetic SkDs and their relatives shown in Fig. [16(e)] are energetically stable on Néel states so that they can be new ingredients of antiferromagnetic spintronics.

V. CONCLUSION

We numerically demonstrated that optical or electron vortex beams can induce a new class of topological magnetic defects in chiral magnets, whose spin texture is given in Fig. [1(b)] for the ferromagnetic case and Fig. [13(b)] [and Fig. [16(e)]] for the antiferromagnetic case.

Supported by quantitative consideration about the time and length scales of chiral magnets, we modeled the effect of vortex beams as heating. The heating realizes spatially modulated temperature proportional to the local intensity of the vortex beams. By solving stochastic Landau-Lifshitz-Gilbert equation for canonical models of chiral ferro- and antiferro- magnets [Eq. (8) and (14)], we found that the ring-shaped profile of the vortex beams, which originates from their non-vanishing intrinsic orbital angular momentum, can be transferred to chiral magnets as topological defects. The new defect, which we named skyrmion duplex, is a bound state of a skyrmion and anti-skyrmion. We note that both ordinary Gaussian beams and vortex beams can create skyrmions, while Gaussian beams without orbital angular momentum cannot be used to make skyrmion duplexes. We confirmed that skyrmion duplexes in chiral ferromagnets are energetically stable in the skyrmion crystal phase. Namely, they have longer lifetime than our calculation periods. For chiral antiferromagnets, they are stable in the broad region of the phase diagram reflecting the degeneracy of Néel states.

We gave optimized way of applying vortex beams to achieve high probability of creating skyrmion duplexes in chiral magnets (Fig. [7] for the ferromagnetic case and Fig. [16] for the antiferromagnetic case). The appropriate wavelength of vortex beams is found to be of the same order of the size of skyrmions. Typically the size of skyrmions in chiral magnets is \(O(1)\sim O(10)\) nano meter so that the proper wavelength is in the extreme ultraviolet region. However,
FIG. 17. Current-driven motion of antiferromagnetic skyrmion duplex in the presence or absence of a magnetic impurity. The impurity has fixed spin direction $\vec{s} = (0, 0, 1)$, which is consistent with the background Néel order, and its position is emphasized as a white dot. We start applying the spin-polarized current $v_x = -0.001$ at $t = 0$. (a) Initial state at $t = 0$ with the impurity. (b) Initial state without the impurity. (c) Final state at $t = 1200$ with the impurity. (d) Final state without the impurity. We see that the antiferromagnetic SkD is pinned by the impurity. The calculations are performed in a periodic system with $300 \times 300$ sites, and here we show $z$-component of staggered spins in a subregion with $200 \times 150$ sites. We place white lines for guide to eyes.

FIG. 18. Collision between an antiferromagnetic skyrmion duplex and a single-site magnetic impurity. (a) Snapshots within the time evolution under the spin-polarized current $v_x = -0.001$ at $t = 0, 1000, 2000,$ and 3000. (b) Snapshots under $v_x = -0.03$ at $t = 0, 40, 80,$ and 120.

recently much larger skyrmions [59–64] are observed in chiral magnetic films in which the optimal wavelength would be lifted to the visible light region.

We also studied basic dynamical properties of skyrmion duplexes in metallic chiral magnets under spin-polarized currents and magnetic impurities. We confirmed that skyrmion duplexes can be pinned by magnetic impurities (Fig. 8 and Fig. 17). In addition, we unveiled that the way of scattering between skyrmion duplexes and impurities changes depending on their relative velocities (see Fig. 9 for chiral ferromagnets and Fig. 18 for chiral anti-ferromagnets). For chiral ferromagnets, transverse drift motion induced by spin transfer torque realizes continuous decomposition of a skyrmion duplex into a pair of a skyrmion and anti-skyrmion (Fig. 10). This offers a unique scheme of creating and controlling relatively long-lived anti-skyrmions. For chiral antiferromagnets, skyrmion duplexes can be driven by much smaller amount of spin-polarized currents than ferromagnetic skyrmion duplexes, and there is no decomposition of skyrmion duplexes under the currents. Therefore, we expect that antiferromagnetic skyrmion duplexes can be new ingredients of antiferromagnetic spintronics and skyrmionics.
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