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Effects of surface anisotropy on magnetic vortex core

Oleksandr V. Pylypovskyi [a,*], Denis D. Sheka [a], Volodymyr P. Kravchuk [b], Yuri Gaididei [b]

[a] Taras Shevchenko National University of Kiev, 01601 Kiev, Ukraine
[b] Institute for Theoretical Physics, 03143 Kiev, Ukraine

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ABSTRACT

The vortex core shape in the three dimensional Heisenberg magnet is essentially influenced by a surface anisotropy. We predict that depending of the surface anisotropy type there appears barrel- or pillow-shaped deformation of the vortex core along the magnet thickness. Our theoretical study is well confirmed by spin–lattice simulations.

1. Introduction

Among different nontrivial magnetization distributions in the nanoscale, magnetic vortices attract a special interest because the vortex configuration can form a ground state in nano- and micron-sized ferromagnets. It takes place when the sample size exceeds the single-domain size due to the competition between an exchange field and a stray one in magnets with small magneto-crystalline anisotropy [1,2]. Nontrivial topological properties of vortices [3] attract interest to their study with perspective application to the high-density magnetic storage devices, nonvolatile magnetic vortex random-access memories [2,4].

In common with stray field effects which favour the vortex configuration, the vortex can form the lowest energy state in magnets with a surface anisotropy [5,6]. Such anisotropy, which always appears in real samples, is originated from the symmetry breaking for the boundary sites of the lattice and can result in the specific uniaxial single-ion anisotropy of different signs [5,7]. In the disk-shaped magnets the edge surface anisotropy can pin the magnetization along the border in the circular, i.e. in the vortex, configuration [5].

Similarity between the effects of the stray field and the surface anisotropy is not casual. Various regimes are known, when the nonlocal dipolar interaction can be approximately reduced to the local effective anisotropy [8–17]. The analysis can be done in some limiting relations between the nanomagnet size $2R$, its thickness L, and exchange length $\ell_{\text{ex}} = \sqrt{A/4\pi M_S^2}$ with A being the exchange constant and M_S being the saturation magnetization. Here we recall the limiting case $L \ll R$ and $\ell_{\text{ex}} \ll R$, where the analytical description [18] shows that the dipolar interaction can be reduced approximately to an on-site inhomogeneous anisotropy energy. In case of disk-shaped particles there appear two effective inhomogeneous anisotropy terms: one is effective anisotropy of face surface charges (easy-plane anisotropy for thin samples and easy-axis anisotropy for thick ones) and another one is effective anisotropy of edge surface charges, which is responsible for the tangential magnetization distribution along the disk edge resulting in clockwise or counterclockwise vortex chirality. Since the first term can change its sign near the disk edge [18] the so-called tailoring vortices [19,20] can also be quantitatively explained by the effective anisotropy model. Nevertheless it should be noted that such a type of vortices cannot be obtained within the simple model we use in this paper.

In this work we study analytically and numerically the influence of the single-ion uniaxial surface anisotropy of different types, easy-surface (ES) and easy-normal (EN), on the three-dimensional (3D) vortex shape for the Heisenberg magnet. We show that the presence of the surface anisotropy breaks the symmetry of magnetization structure in the axial $\hat{\boldsymbol{z}}$–direction, which naturally leads to $\hat{\boldsymbol{z}}$–dependence of the vortex core width: there appears the barrel- and the pillow-deformation of the core for the ES and EN anisotropies, respectively.

The paper is organized as follows: In Section 2 we introduce a mathematical model of classical Heisenberg ferromagnet with account of the surface anisotropy. We use the variational approach to describe the structure of magnetic vortex in Section 3. Our analytical results are verified in Section 4 by spin–lattice simulations. We discuss in Section 5 how our predictions about the influence of the surface anisotropy on the vortex structure can be applied to magnetic particles, where the stray field effects can be essential. The mathematical details for deriving an energy functional is placed in Appendix A.

* Corresponding author.
E-mail address: engraver@univ.net.ua (O.V. Pylypovskyi).

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2. The model

The model we consider is a ferromagnetic system, described by the classical anisotropic Heisenberg Hamiltonian

$$\mathcal{H} = -J\mathcal{S}^2 \sum_{(\boldsymbol{n},\boldsymbol{\delta})} \boldsymbol{m_n}\cdot\boldsymbol{m}_{\boldsymbol{n}+\boldsymbol{\delta}} + \mathcal{H}^{\text{an}}, \tag{1a}$$

where $J>0$ is the exchange integral, $\mathcal{S}$ is the length of classical spin, $\boldsymbol{m_n}$ is the normalized magnetic moment on a 3D site position $\boldsymbol{n}$, the 3D index $\boldsymbol{\delta}$ runs over the nearest neighbours, and $\mathcal{H}^{\text{an}}$ is the anisotropy part of the Hamiltonian. We take into account the bulk on-site anisotropy with the constant $K>0$ (easy-plane anisotropy) and the surface one with the surface anisotropy constant K_s [7,21]:

$$\mathcal{H}^{\text{an}} = \frac{K\mathcal{S}^2}{2}\sum_{\boldsymbol{n}}(\boldsymbol{m_n}\cdot\hat{\boldsymbol{z}})^2 - \frac{K_s\mathcal{S}^2}{2}\sum_{(\boldsymbol{l},\boldsymbol{\delta})}(\boldsymbol{m_l}\cdot\boldsymbol{u_{l\delta}})^2. \tag{1b}$$

Here the unit vector $\hat{\boldsymbol{z}}$ is the cylinder axis and the last term describes the Néel surface anisotropy with the unit vector $\boldsymbol{u_{l\delta}}$ connecting the magnetic moment $\boldsymbol{m_l}$ from the surface site $\boldsymbol{l}$ to its nearest neighbour $\boldsymbol{\delta}$.

The continuum description of the system is based on smoothing of the lattice model using the normalized magnetization

$$\begin{aligned}\boldsymbol{m}(\boldsymbol{r},t) &= a^3\sum_{\boldsymbol{n}}\boldsymbol{m_n}\delta(\boldsymbol{r}-\boldsymbol{r_n})\\ &= (\sin\theta\cos\phi,\ \sin\theta\sin\phi,\ \cos\theta),\end{aligned} \tag{2}$$

where $\theta=\theta(\boldsymbol{r},t)$, $\phi=\phi(\boldsymbol{r},t)$, the parameter a being the lattice constant, and $\delta(\boldsymbol{r})$ being the Dirac δ-function.

The total energy, the continuum version of the Hamiltonian (1), normalized by $K\mathcal{S}^2/a^3$ has the following form:

$$\begin{aligned}\mathcal{E} &\equiv \frac{E}{K\mathcal{S}^2/a^3} = \mathcal{E}_v + \mathcal{E}_s,\\ \mathcal{E}_v &= \frac{1}{2}\int dV\,[-\ell^2\boldsymbol{m}\cdot\nabla^2\boldsymbol{m} + (\boldsymbol{m}\cdot\hat{z})^2],\\ \mathcal{E}_s &= \frac{\varkappa a}{2}\int dS\,(\boldsymbol{m}\cdot\boldsymbol{n}_s)^2\end{aligned} \tag{3}$$

with $\ell = a\sqrt{J/K}$ being the magnetic length which is a natural scale in the model where only exchange and anisotropy energies are taken into account.[1] The last term $\mathcal{E}_s$ is the transverse surface anisotropy, the continuum analogue of Néel surface anisotropy with $\boldsymbol{n}_s$ being the normal to the surface and the parameter $\varkappa = K_s/K$ being the surface anisotropy rate. In the further study we consider the cases of both ES anisotropy when $\varkappa>0$ and the EN one when $\varkappa<0$.

The equilibrium magnetization structure can be found by varying the energy functional (3), which results in the following boundary-value problem [1,22]:

$$\boldsymbol{m}\times[\ell^2\nabla^2\boldsymbol{m} - (\boldsymbol{m}\cdot\hat{\boldsymbol{z}})\hat{\boldsymbol{z}}] = 0, \tag{4a}$$

$$\ell^2\frac{\partial\boldsymbol{m}}{\partial\boldsymbol{n}_s}\bigg|_S = \varkappa a(\boldsymbol{m}\cdot\boldsymbol{n}_s)[(\boldsymbol{m}\cdot\boldsymbol{n}_s)\boldsymbol{m} - \boldsymbol{n}_s]|_S. \tag{4b}$$

The absence of nonlocal dipolar interaction allows one to avoid integrodifferential equations here [23,24]. One can see that the presence of the surface anisotropy changes the symmetry of boundary conditions, leading to the Robin boundary conditions instead of the Neumann ones [25]. As a result the symmetry breaking the magnetization structure becomes $\hat{z}-$dependent. In particular, we will see that the vortex core width takes different values in a volume and on the surface.

[1] Typical characteristic lengths are of 10 nm like exchange lengths $\ell_{\text{ex}} = 5.1$ nm for Permalloy, 7.6 nm for Nickel and magnetic length 4.7 for Cobalt [2]. Magnetostatics can be reduced to the effective surface anisotropy in thin films and in this case the vortex core size is a corresponding effective magnetic length.

3. Vortex core structure: analytics

Let us consider the disk-shape sample with the radius R and the thickness L. The volume contribution to the energy functional (3) reads

$$\mathcal{E}_v = \frac{1}{2}\int dV\,\left\{\ell^2[(\nabla\theta)^2 + \sin^2\theta(\nabla\phi)^2] + \cos^2\theta\right\}. \tag{5a}$$

The surface energy term $\mathcal{E}_s = \mathcal{E}_s^{\text{face}(+)} + \mathcal{E}_s^{\text{face}(-)} + \mathcal{E}_s^{\text{edge}}$,

$$\begin{aligned}\mathcal{E}_s^{\text{face}(\pm)} &= \frac{\varkappa a}{2}\int dS^{\text{face}(\pm)}\,\cos^2\theta|_{z=\pm L/2},\\ \mathcal{E}_s^{\text{edge}} &= \frac{\varkappa a}{2}\int dS^{\text{edge}}\,\sin^2\theta\,\cos^2(\phi-\chi)|_{\rho=R},\end{aligned} \tag{5b}$$

where (ρ,χ,z) are the cylinder coordinates.

In terms of the angular variables the boundary-value problem (4) for the disk-shaped sample has the following form:

$$\nabla^2\theta - \frac{1}{2}\sin 2\theta\left[(\nabla\phi)^2 - \frac{1}{\ell^2}\right] = 0, \tag{6a}$$

$$\nabla\cdot(\sin^2\theta\nabla\phi) = 0, \tag{6b}$$

$$\pm\ell^2\frac{\partial\theta}{\partial z} - \frac{\varkappa a}{2}\sin 2\theta\bigg|_{z=\pm L/2} = 0,\quad \frac{\partial\phi}{\partial z}\bigg|_{z=\pm L/2} = 0, \tag{6c}$$

$$\ell^2\frac{\partial\theta}{\partial\rho} + \frac{\varkappa a}{2}\sin 2\theta\,\cos^2(\phi-\chi)\bigg|_{\rho=R} = 0, \tag{6d}$$

$$\ell^2\frac{\partial\phi}{\partial\rho} - \frac{\varkappa a}{2}\sin 2(\phi-\chi)\bigg|_{\rho=R} = 0. \tag{6e}$$

The form of boundary conditions determines possible minimizers. One can see that the boundary-value problem (6) has the vortex-like stationary solution with

$$\phi = \chi + \varphi_0. \tag{7a}$$

To satisfy the boundary condition (6e), the value of the constant $\varphi_0 = \pm\pi/2$ for $\varkappa>0$ (ES magnets) and $\varphi_0 = 0,\pi$ for $\varkappa<0$ (EN magnets).

The simplified version of the boundary-value problem (6) with $\theta=\pi/2$ was considered in Refs. [5,13,17]: Planar vortices with $\cos\theta = 0$ and $\phi = \chi + \varphi_0$ were shown to be metastable states in the disk-shaped system.

Below we discuss the 3D boundary-value problem (6). In this case the nonplanar vortex with z-dependence of the polar angle appears as follows:

$$\theta = \theta(\rho, z). \tag{7b}$$

The typical scale of the θ-distribution is determined by the magnetic length ℓ. Supposing that $\ell \ll R$, we can replace the boundary condition (6d) by

$$\frac{\partial\theta}{\partial\rho}\bigg|_{\rho=R\to\infty} = 0,\qquad \cos\theta|_{\rho=R\to\infty} = 0. \tag{8}$$

The problem (6) is the nonlinear boundary-value problem for the partial differential equation for the function (7b). To simplify the analysis we use the variational approach with Ansatz-function:

$$\cos\theta(\rho,z) = f\left(\frac{\rho}{w(z)\ell}\right),\quad f(x) = \exp\left(-\frac{x^2}{2}\right). \tag{9}$$

This function is the generalization of the well-known Feldt-keller Ansatz [1,26], originally used to describe the structure of the vortex in thin films. However in contrast to [1] our reduced vortex core function $w(z)$ is a variational function.

Using Ansatz (9) one can write down the energy in the form $\mathcal{E} = \mathcal{E}_0 + \pi\ell^3\sqrt{\zeta(3)}\tilde{\mathcal{E}}[w]$, where the first term $\mathcal{E}_0$ is independent of the z coordinate, the second term $\tilde{\mathcal{E}}[w]$ contains terms both due