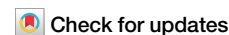


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Antiferromagnetic domain wall dynamics in rotating magnetic fields engineered by bending and twisting

Yelyzaveta A. Borysenko^{1,2,3}, Kostiantyn V. Yershov^{2,4} & Denis D. Sheka¹ ✉

We study domain wall motion in an intrinsically achiral helix-shaped antiferromagnet driven by a circularly polarized magnetic field, using the collective-variable approach and spin-lattice simulations. We show that this problem can be reduced to an effective dynamics of the domain wall in a rectilinear chiral biaxial antiferromagnet in the three rotating spatially-dependent fields. Two distinct modes of motion are identified: (i) rigid domain wall motion at low frequencies, characterized by constant velocity and a fixed domain wall phase; and (ii) oscillatory motion at high frequencies, accompanied by precessional dynamics of the domain wall. In the limit case of zero curvature, the domain wall becomes immobile; however, its phase remains fully synchronized with the external field, with the direction of phase rotation determined by the geometrical chirality of the helix. We predict that the domain wall velocity in the rigid mode can be efficiently tuned by both the field frequency and geometric manipulations via bending and twisting. These findings suggest that 3D curvilinear antiferromagnets such as helical structures may serve as a platform for the future high speed antiferromagnetic curvilinear spintronics applications.

Antiferromagnetic (AFM) spintronics is a fast developing topic in spintronics that attracts high interest due to high stability under magnetic fields, terahertz operation, and memory without stray fields^{1–3}. One of the obstacles compared to traditional ferromagnetic spintronics is the dramatically reduced thermal stability of AFM nanostructures compared to their FM counterparts⁴. To overcome the problem, one can encode information by a topologically protected object like a domain wall. Today AFM domain walls can be efficiently driven by different stimuli such as spin wave excitations^{5–7}, inhomogeneous static magnetic fields⁸, rotating magnetic fields^{9–11}, spin-orbit torques^{12–17}, spin-transfer torques^{18–21}, strains²², and thermal gradients^{23–25}. AFM materials with chiral interactions are of special interest due to the perspective of applications in chiral spintronics²⁶. One of the well-known chiral interaction appears in systems with broken inversion symmetry: this is the case of Dzyaloshinskii–Moriya interaction (DMI), which favors canted textures, spin spirals, and chiral skyrmions^{27,28}. Recent studies of domain walls (DW) in AFM materials show that DMI affects both statics and dynamics, including DW tilting, reduction of the wall width, and selectivity of chirality²⁹. DWs in AFM driven by magnons can move an order of magnitude faster in the presence of DMI³⁰ with polarization effect³¹. An efficient way to drive DWs in AFMs with DMI can be realized by rotating magnetic fields with speeds much faster than in ferromagnets^{9,10,32}.

A number of efficient ways of manipulation by magnetization texture in AFMs can be realized in curved magnetic architectures. By exploring the mutual interplay between curved geometry, nontrivial topology of the underlying substrate, and spin texture, curvilinear magnetism offers a promising playground for control and manipulation of magnetic response by tailoring geometrical properties³³. The topic of curvilinear magnetism is well established for ferromagnetic materials, while curvilinear antiferromagnetism is just beginning to develop³⁴. In contrast to ferromagnets, achiral AFM spin chains possess homogeneous DMI and emergent weak ferromagnetism³⁵. Inhomogeneous DMI in achiral spin chains results in the spin-flop transition in AFM ring chains being of first or second order, depending on the ring curvature³⁶. Besides, an intrinsically achiral curvilinear AFM spin chain behaves as a chiral helimagnet with geometrically tunable DMI³⁷. Curvature drastically modifies the dynamic response of AFMs. In particular, geometry-governed inhomogeneous DMI in helical AFM is responsible for the hybridization of spin wave modes³⁷. Polarized spin waves can be used in a spin-wave rotator and a spin-wave interferometer based on magnetic circular helices, which offer new ways in constructing magnonic logic devices³⁸.

Being crucial for the applications, geometry-governed effects of DW dynamics in curved systems, including spin chains, wires, stripes, and

¹Taras Shevchenko National University of Kyiv, Kyiv, Ukraine. ²Leibniz-Institut für Festkörper- und Werkstofforschung, IFW Dresden, Dresden, Germany.

³Department of Physics, University of Konstanz, Konstanz, Germany. ⁴Bogolyubov Institute for Theoretical Physics of the National Academy of Sciences of Ukraine, Kyiv, Ukraine. ✉e-mail: sheka@knu.ua

ribbons are well established for the ferromagnets only, see for the review³⁹ with applications for curvilinear spintronics⁴⁰. Recently, DW dynamics was studied in thin planar AFM wires⁴¹: the DW is pinned at the curvature extremum for the case of localized bend, the DW can be efficiently driven by the curvature gradient, which results in its automotion⁴¹. Rotating magnetic fields can act as an effective torque which can assist or drive magnetization switching⁴². In straight AFM spin chains with intrinsic DMI, DWs can be efficiently driven by a rotating magnetic field. Such a motion exhibits a Walker-breakdown behavior with a critical frequency ω_c proportional to DMI constant. The threshold value ω_c separates two regimes^{9,11,32}: (i) A phase-locking regime with rigid DW motion is realized at low frequencies, when $\omega < \omega_c$ it correspond to the steady state motion of DW with a speed v proportional to the field frequency ω . The high frequency ($\omega > \omega_c$) phase-unlocking regime is accompanied by DW oscillations.

The goal of this work is to achieve geometric control of DW dynamics in intrinsically achiral AFM assisted by a rotating magnetic field. We illustrate the influence of the curved geometry by an example of a helix-shaped AFM spin chain. Using a methodology of curvilinear antiferromagnetism, reviewed in³⁴, we study how bending and twisting act on the DW motion. It is well-known that exchange interaction in curved geometry becomes a source of two emergent interactions: in AFM helices there appear geometry-governed exchange-driven DMI and biaxial anisotropy³⁷. If a DW is driven in a helix by a rotating magnetic field, one can expect phase-locking regimes at low frequencies and phase-unlocking regimes for high frequencies, similar to the case of intrinsic DMI^{9,11,32}. Surprisingly, the DW dynamics in helices demonstrate an opposite picture: the low-frequency rigid motion is not characterized by phase synchronization, while the high-frequency regime is characterized by phase oscillation motion.

Results

Model of curvilinear antiferromagnet

We start with an intrinsically achiral one-dimensional antiferromagnetic spin chain, in which the magnetic moments μ_i are arranged along a helix-shaped curve $\gamma(s)$, with $s = \delta_a i$ denoting the arc-length coordinate, where δ_a is the lattice constant. The helix-shaped curve is defined as $\gamma(s) = R[\cos(s/s_0)\hat{x} + \sin(s/s_0)\hat{y}] + \frac{CP}{2\pi}(s/s_0)\hat{z}$ in Cartesian basis, where R, P are radius and pitch of the helix, respectively [see Fig. 1(a)], curvature and torsion are defined as $\kappa = R/s_0^2$, $\tau = CP/(2\pi s_0^2)$ where $s_0 = \sqrt{R^2 + P^2}/(4\pi^2)$. Parameter $C = \pm 1$ is the helix chirality (left- or right-handed helix). In magnetic interactions, we consider (i) nearest-neighbor AFM exchange $\mathcal{H}_X = \frac{1}{2} \mathcal{J} S^2 \sum_i \mu_i \cdot \mu_{i+1}$ with exchange integral $\mathcal{J} > 0$ and S being the spin length. (ii) Easy-axial anisotropy $\mathcal{H}_A = -\frac{1}{2} \mathcal{K} S^2 \sum_i (\mu_i \cdot \mathbf{e}_T^i)^2$ with $\mathcal{K} > 0$ being easy-axial anisotropy strength. The anisotropy axis $\mathbf{e}_T^i$ is spatially dependent and is tangential to the curve γ . (iii) The interaction with an external magnetic field $\mathcal{H}_Z = -g\mu_B S \sum_i \mu_i \cdot \mathbf{H}$ with g being Landé factor and μ_B being Bohr magneton. We consider static and time-dependent magnetic fields, the latter rotates in the plane, perpendicular to the symmetry axis of the helix ($\hat{z}$) with constant frequency ω , see Fig. 1(a).

It is important to mention that the considered Hamiltonian does not include intrinsic chiral interaction (such as bulk/interfacial DMI); consequently, the magnetic system is intrinsically achiral in the sense that left- and right-handed helices are symmetry-related and energetically equivalent in the absence of a field. This follows from the invariance of the exchange-anisotropy Hamiltonian under (i) sublattice interchange and (ii) reversal of the curve parametrization $s \rightarrow -s$, which reverses the sign of the torsion and, hence, the geometric handedness.

In the following, we utilize the continuum description in which magnetization of each sublattice is represented by a continuous vector function $\mathbf{M}_\alpha(s) = M_s \mu^\alpha$ of the constant absolute value $|\mathbf{M}_\alpha| = M_s$. Here $M_s = g\mu_B S/(2\delta_a)$ is the saturation magnetization of one sublattice. It is instructive to introduce the dimensionless vector of total magnetization $\mathbf{m} = (\mathbf{M}_1 + \mathbf{M}_2)/(2M_s)$ and Néel vector $\mathbf{n} = (\mathbf{M}_1 - \mathbf{M}_2)/(2M_s)$, where $\mathbf{n} \cdot \mathbf{m} = 0$ and $\mathbf{n}^2 + \mathbf{m}^2 = 1$. In the limit $|\mathbf{m}| \ll 1$, dynamics of the Néel vector can be described within

the Lagrange formalism^{43,44}. The Lagrangian $L = \int_{-\infty}^{+\infty} \mathcal{L} ds$ and Rayleigh dissipative function $R = \int_{-\infty}^{+\infty} \mathcal{R} ds$ are

$$\begin{aligned} \mathcal{L} &= \frac{M_s^2}{\gamma_0 \Lambda} [\partial_t \mathbf{n} + \mathbf{n} \times \gamma_0 \mathbf{H}]^2 - \mathcal{E}, \quad \mathcal{R} = \frac{M_s^2}{\gamma_0 \Lambda} \frac{\eta}{\zeta} (\partial_t \mathbf{n})^2, \\ \mathcal{E} &= A (\partial_s \mathbf{n})^2 + K [1 - (\mathbf{n} \cdot \mathbf{e}_T)^2], \end{aligned} \quad (1)$$

where γ_0 is the gyromagnetic ratio, η is a damping coefficient, $\Lambda = 2\mathcal{J}S^2/(\delta_a)$ is the uniform exchange constant, $A = \mathcal{J}S^2\delta_a$ is the exchange stiffness constant, $K = \mathcal{K}S^2/(2\delta_a)$ is the anisotropy constant. The model (1) is valid for a strong exchange field, i.e. $H_A/H_X = K/\Lambda = \zeta^2 \ll 1$.

In the following, it is convenient to consider the energy E from (1) in the curvilinear frame of reference. For the case of zero external magnetic field effective geometry-induced interactions cause a tilt of the equilibrium Néel vector from the easy-axial, tangential $\mathbf{e}_T$, direction [see Fig. 1(b)] by an angle $\psi \approx \kappa\sigma^{45}$, where $\kappa = \ell\kappa$ and $\sigma = \ell\tau$ are dimensionless curvature and torsion, respectively, with $\ell = \sqrt{A/K}$ being the magnetic length which defines the length scale of the system. This tilt suggests to consider the system in the rotated reference ψ -frame $\{\mathbf{e}_1, \mathbf{e}_2, \mathbf{e}_3\}$ ⁴⁶, which diagonalizes the effective anisotropy energy density and align Néel vector along the curvilinear basis. The energy density of the system normalized by K in the rotated ψ -frame reads³⁷

$$\begin{aligned} \tilde{\mathcal{E}} &= n'_k n'_k - \mathcal{H}_1 n_1^2 + \mathcal{H}_2 n_2^2 \\ &+ \mathcal{D}_1 (n_2 n'_3 - n_3 n'_2) + \mathcal{D}_2 (n_1 n'_2 - n_2 n'_1), \end{aligned} \quad (2)$$

where prime denotes to the derivative with respect to the dimensionless coordinate $\xi = s/\ell$, and n_k are components of the Néel vector $\mathbf{n} = (\cos \theta, \sin \theta \cos \phi, \sin \theta \sin \phi)$ in the ψ -frame. The first term in energy (2) corresponds to the common inhomogeneous exchange interaction in a rectilinear system. The second and third terms correspond to the effective easy- and hard-axial anisotropy terms, respectively, with $\mathcal{H}_1 \approx 1 - \kappa^2 + \sigma^2$ and $\mathcal{H}_2 \approx \kappa^2$ being the corresponding anisotropy coefficients. The last two terms include Lifshitz-like invariants which describe a geometry-induced DMI with $\mathcal{D}_1 \approx 2\sigma$ and $\mathcal{D}_2 \approx 2\kappa$ being geometry-induced DMI coefficients. These Lifshitz invariants are geometry-induced and remain symmetry-compatible with intrinsic achirality.

Field-induced DW dynamics

We consider DW dynamics in helix-shaped geometry driven by the rotating magnetic field in the xy -plane $\mathbf{h} = h(\hat{x} \cos \omega \bar{t} + \hat{y} \sin \omega \bar{t})$ with frequency $\omega = \omega/\omega_0$ and dimensionless time $\bar{t} = \omega_0 t$. Here $\omega_0 = \gamma_0 \sqrt{\Lambda K}/M_s$ is the frequency of uniform AFM resonance, $h = H/H_{sf}$ is dimensionless magnetic field with $H_{sf} = \sqrt{\Lambda K}/M_s$ being the spin-flop field. In the ψ -frame, this field can be presented as a sum of three fields that rotate in different planes

$$\begin{aligned} \mathbf{h} &= h(\mathbf{h}_{23} + \mathbf{h}_{13} + \mathbf{h}_{12}), \\ \mathbf{h}_{23} &= -C_\xi \left[\mathbf{e}_2 \cos \omega \bar{t} + \mathbf{e}_3 \frac{\kappa \sin \psi + \sigma \cos \psi}{\sqrt{\sigma^2 + \kappa^2}} \sin \omega \bar{t} \right], \\ \mathbf{h}_{12} &= -S_\xi \left[\mathbf{e}_1 \frac{\kappa \cos \psi - \sigma \sin \psi}{\sqrt{\sigma^2 + \kappa^2}} \cos \omega \bar{t} + \mathbf{e}_2 \sin \omega \bar{t} \right], \\ \mathbf{h}_{13} &= \mathbf{e}_1 \frac{\kappa \cos \psi - \sigma \sin \psi}{\sqrt{\sigma^2 + \kappa^2}} C_\xi \cos \omega \bar{t} \\ &+ \mathbf{e}_3 \frac{\kappa \sin \psi + \sigma \cos \psi}{\sqrt{\sigma^2 + \kappa^2}} S_\xi \sin \omega \bar{t}, \end{aligned} \quad (3)$$

where $C_\xi = \cos(\xi\sqrt{\sigma^2 + \kappa^2})$ and $S_\xi = \sin(\xi\sqrt{\sigma^2 + \kappa^2})$. Considering model (2) together with the external field given in (3), the problem of DW dynamics in an achiral uniaxial helix-shaped system in an external magnetic field rotating in the xy -plane can be reformulated as an effective problem of DW dynamics in a chiral biaxial “rectilinear” system in three spatially dependent fields, each rotating in a different plane.

The effective equations of motion for the DW are the Euler-Lagrange-Rayleigh equations with the effective Lagrangian and dissipation functions

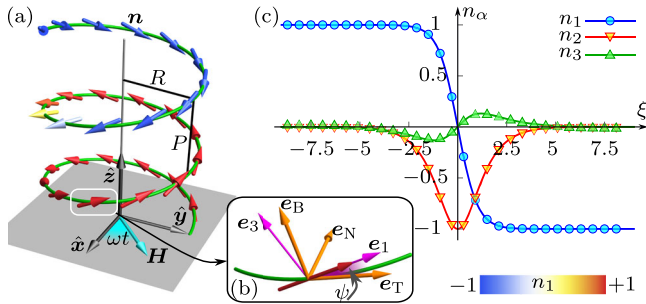


Fig. 1 | Domain wall in AFM helix. **a** Schematic representation of the AFM DW in a helix-shaped geometry. External magnetic field H is applied in the xy plane, rotating with the frequency ω . **b** TNB and rotated ψ -reference frames and the Néel vector in equilibrium quasitangential state being tilted from the axis e_T by the angle ψ . **c** DW profile: solid lines represent model (8), symbols correspond to the results of spin-lattice simulations.

normalized by $E_0 = \sqrt{AK}$ for the DW ansatz (8) are

$$\mathcal{L} = \sum_{v=0}^2 h^v \mathcal{L}_v - \mathcal{E}^{\text{DW}}, \quad \mathcal{R} = \tilde{\eta} \mathcal{L}_0, \quad (4)$$

where $\tilde{\eta} = \eta/\zeta$ is the normalized damping coefficient. Due to the complex structure of the external magnetic field (3), the explicit form of $\mathcal{L}_v$ terms is bulk for presentation and can be found in the Supplemental materials (SM). The DW energy $\mathcal{E}^{\text{DW}}$ is presented in Methods.

The general form of the equations of motion reads

$$\begin{aligned} \widehat{\mathbf{M}} \cdot \begin{bmatrix} \ddot{\mathbf{q}} \\ \dot{\Phi} \end{bmatrix} + (\widehat{\mathbf{R}} + \tilde{\eta} \widehat{\mathbf{M}}) \cdot \begin{bmatrix} \dot{\mathbf{q}} \\ \dot{\Phi} \end{bmatrix} &= \mathbf{F}, \quad \mathbf{F} = \begin{bmatrix} F_q \\ F_\Phi \end{bmatrix}, \\ \widehat{\mathbf{M}} &= 4\Delta \begin{bmatrix} -\frac{1+a^2\Delta^2}{\Delta^2} & a \\ a & -1 \end{bmatrix}, \quad \widehat{\mathbf{R}} = \begin{bmatrix} R_q & R_{q\Phi} \\ -R_{q\Phi} & 0 \end{bmatrix}, \end{aligned} \quad (5)$$

where overdot indicates the derivatives with respect to dimensionless time, $\widehat{\mathbf{M}}$ is an effective DW mass matrix, $\widehat{\mathbf{R}}$ field-induced effective damping matrix, and $\mathbf{F}$ field-induced force. Parameters R_i and F_i are functions of DW position $q(t)$, phase $\Phi(t)$, geometrical parameters κ , σ , and external field h , ω .

Considering the curvature and torsion as small parameters within a low-frequency regime, we can get a traveling wave solution, a *rigid* DW motion mode, with $\dot{\mathbf{q}} = v$ being a constant velocity and $\varphi = \Phi - \Phi_0$ being a small deviation of the phase from the equilibrium state, $|\varphi| \ll 1$, see Fig. 2. For this we consider (5) in a harmonic approximation, where we save terms which are quadratic with respect to small parameters κ , σ , φ , and ω in the equations of motion (5) and integrated it with respect to time $\bar{t} \in [0, 2\pi/\omega]$. As a result, we obtain the DW velocity

$$v = v_0 + \delta v, \quad v_0 = \omega / \sqrt{\sigma^2 + \kappa^2}, \quad (6)$$

and DW phase deviation

$$\varphi = \varphi_0 + \delta\varphi, \quad \varphi_0 = 2h\omega\kappa \frac{(\pi^2 - 1)\sqrt{\kappa^2 + \sigma^2}}{(2\kappa^2 + \sigma^2)^2}, \quad (7)$$

where pairs $\{v_0, \varphi_0\}$ and $\{\delta v, \delta\varphi\}$ are considered as asymptotic DW velocity/phase deviation and its small corrections for the low damping case, respectively, which is pronounced only at small curvatures, see dashed lines in Fig. 3(a). The corrections for velocity $\delta v \propto \omega^2$ and phase $\delta\varphi \propto \omega^2 \tilde{\eta}^2$ are of second order of magnitude with respect to frequency, for details see (SM). It should be noted that the asymptotic velocity given by (6) is valid only for finite curvature ($\kappa > 0$). In the limiting case of zero curvature but finite torsion, the system exhibits a qualitatively different regime of DW motion, which is discussed below.

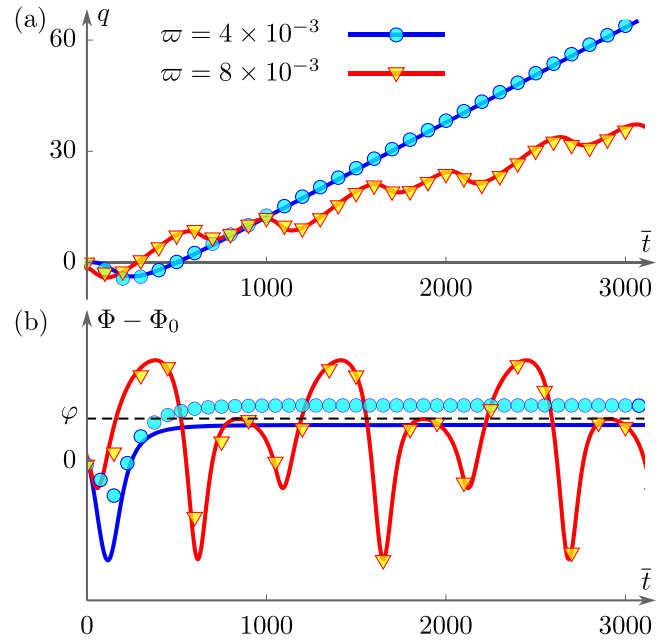


Fig. 2 | Temporal evolution of domain walls in AFM helices. Time dependencies of DW position **(a)** and phase **(b)** in rigid (blue lines) and oscillating (red lines) motion modes for field frequencies $\omega = 4 \times 10^{-3}$ and $\omega = 8 \times 10^{-3}$, respectively. Solid lines and symbols correspond to numerical solution of set (5) and spin-lattice simulations, respectively. Other parameters: $\kappa = 0.12$, $\sigma = 0.1$, $h = 0.2$, $\zeta = 5 \times 10^{-2}$, and damping $\eta = 10^{-2}$. Horizontal dashed black line in **(b)** corresponds to the φ defined in (7).

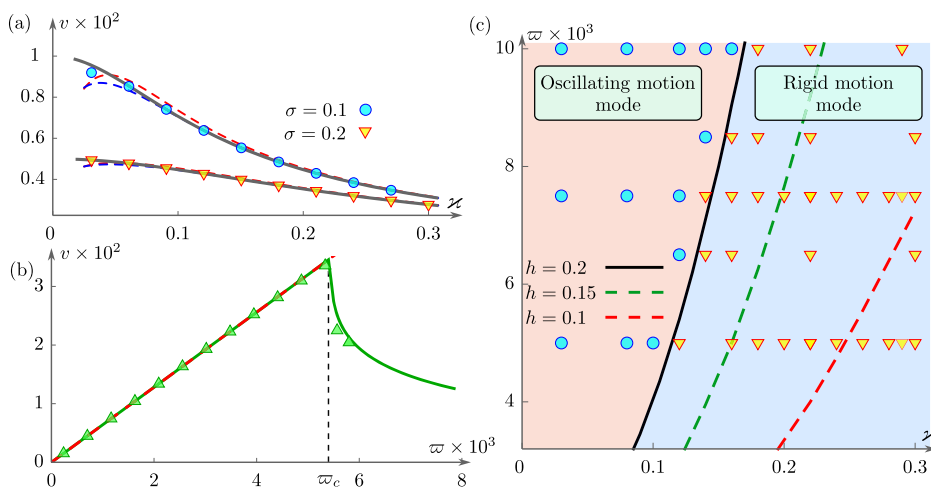
DW asymptotic velocity (6) is proportional to the frequency of the field ω that agrees with results by Pan et al.⁹. However, in our case, we observe the regime of the fixed phase with $\Phi = 0$ for the traveling wave solution instead of the phase-locking regime with $\Phi = \omega$, observed in ref. 9. The origin of the fixed phase regime is emergent geometry-driven DMI. The torsion and curvature result in the appearance of phase slope $a_0 \approx -\sigma$ and phase selectivity $\cos \Phi_0 = -\text{sign}(p\kappa)$, respectively. The last effect results in the fixed phase, see Eq. (S12) of (SM) during the dynamics, because any deviation of phase from its equilibrium value increases the total DW energy, see Eq. (11). Therefore, the transition to the *synchronized motion*⁹ regime becomes suppressed.

The traveling wave solution (6) exists for the frequencies $\omega < \omega_c$, where ω_c is a critical frequency value which separates the rigid and oscillating motion modes, see Figs. 2 and 3(b). We found the oscillating motion mode in simulations and numerical solutions of set (5): we were unable to obtain an analytical characterization of this regime. The oscillation mode is similar to the Walker breakdown in ferromagnets, where DW velocity rapidly drops down and DW starts to precess, see Figs. 2 and 3(b). The amplitude of external magnetic field does not influence the DW velocity, however, h stabilizes the rigid motion mode (i.e., it defines the phase deviation φ): for smaller h the critical frequency ω_c that separates oscillating and rigid motion modes, is shifted to the larger curvatures, see dashed lines for $h = 0.1$ and $h = 0.15$ in Fig. 3(c).

Although we consider an antiferromagnetic spin-chain with easy-axis anisotropy, we estimate the value of the DW velocity using parameters for NiO ⁴⁷ and KFeS_2 ⁴⁸ materials with magnon velocities $v_{\text{NiO}} = \omega_0^{\text{NiO}} \ell^{\text{NiO}} \approx 30.6$ km/s and $v_{\text{KFeS}_2} = \omega_0^{\text{KFeS}_2} \ell^{\text{KFeS}_2} \approx 125.3$ km/s, respectively. Using these parameters with geometrical parameters $\kappa = 0.12$ and $\sigma = 0.1$, one can estimate the maximal DW velocities for a rigid DW motion (i.e., for *traveling wave solution*) as $v_0^{\text{NiO}} \approx 1.06$ km/s and $v_0^{\text{KFeS}_2} \approx 4.33$ km/s. These values are higher than in the rectilinear system⁹, but comparable to the cases where DW motion is induced by spin-orbit torques^{12,49} and coherent spin waves⁵⁰.

For one-dimensional spin chains, one can consider single-chain magnets, *molecular chains*, in which metal atoms are arranged along a helical backbone and surrounded by radicals. Metal atoms and radicals

Fig. 3 | Two types of domain wall motion. **a** DW velocity as a function of helix curvature in the rigid motion mode for helices with $\sigma = 0.1$, $\sigma = 0.2$, and external field with $\omega = 10^{-3}$ and $h = 0.2$. Solid lines represent asymptotic velocity v_0 , dashed lines correspond to (6) for $\eta = 10^{-2}$ (blue line) and $\eta = 10^{-3}$ (red line). **b** Averaged DW velocity as a function of external magnetic field frequency for the helix with $\kappa = 0.12$, $\sigma = 0.1$, and field $h = 0.2$. The red dashed line corresponds to the asymptotic DW velocity v_0 defined in (6). Vertical dashed line defines the critical frequency ω_c that separates rigid and oscillating motion modes. **c** ω - κ phase diagram of the DW motion regimes for helices with $\sigma = 0.1$ and fields with $h = 0.2$, solid line corresponds to the boundary between the states obtained for $h = 0.2$, dashed lines correspond to the boundaries for different external field amplitudes, $h = 0.15$ and $h = 0.1$. In all panels, symbols represent data obtained from spin-lattice simulations, lines correspond to values obtained from numerical solutions of (5), and we consider $\eta = 10^{-2}$ and $\zeta = 5 \times 10^{-2}$.



possess spins and couple to form ferro-, ferri-, or antiferromagnetic ordering within the chain^{51–53}. By choosing different metal ions and radicals, the magnetic anisotropy strength^{52,54} and the exchange coupling can be tuned (the latter can also be controlled with electric fields in certain compounds⁵⁵). For a cobalt(II) system, $\text{Co}(\text{hfac})_2(\text{EtO-NapNIT})$ ⁵⁴, helices with geometrical parameters $\kappa = 0.63$ and $\sigma = 0.13$ can be synthesized. In such a system, $\ell^{\text{Co-chain}} \approx 1.5$ nm and $\omega_0^{\text{Co-chain}} \approx 17$ THz, enabling DW motion with estimated velocities up to $v_0^{\text{Co-chain}} \sim 1$ km/s. By varying the radicals in cobalt-based chains, both geometric and magnetic properties can be modified while preserving the helical architecture and strong cobalt-driven anisotropy^{53,56,57}. The estimated DW velocity for a single-chain magnet is about two orders of magnitude higher than velocities achieved in rectilinear systems driven by a rotating field⁹, by spin waves⁶, or by spin-orbit torques¹².

Limit case of zero curvature

Next, we consider the DW dynamics in the limit case of zero curvature and a finite torsion σ , which corresponds to the straight spin chain with rotating anisotropy axis. In this case, the external magnetic field (3) reduces to the simple form $\mathbf{h} = h[\mathbf{e}_3 \sin \xi \sigma - \mathbf{e}_2 \cos \xi \sigma] \cos \omega t - h[\mathbf{e}_3 \sin(\sigma) \cos \xi \sigma - \mathbf{e}_2 \sin \xi |\sigma|] \sin \omega t$. Similar to the previous case, here we reduce our problem to the problem of DW dynamics in chiral uniaxial system in the relating spatially-dependent field. Spatial dependence in the external field appears due to nonzero torsion, which is a source of the rotation of the reference frame. The effective Lagrangian function for DW in this case becomes simpler, while the Rayleigh function remains the same, for details see Methods and (SM).

In the limit case of zero curvature, we observe synchronization between the DW phase and the external magnetic field. However, despite this phase synchronization, the DW exhibits zero velocity and remains immobile, i.e., $\dot{q} = 0$. The DW phase rotates with the same frequency as the external field, $\Phi = C\omega$, similar to the rectilinear system⁹. Importantly, (i) the direction of the phase rotation is determined by the sign of the torsion ($C = \text{sign}(\sigma)$), which allows for situations where the DW phase and the external field rotate in opposite directions; (ii) the presence of finite curvature is crucial to have a non-zero DW velocity.

Discussion

Our findings open perspectives for applications in antiferromagnetic spintronics, where DWs act as robust information carriers. The ability to drive and control AFM DW dynamics in helical geometries using rotating magnetic fields provides a pathway toward racetrack-type memories and logic devices that exploit the absence of stray fields and terahertz operation

speeds^{1,2}. The geometry-induced effective DMI enables chiral functionalities even in achiral systems, which is promising for chiral spintronics^{26,29,30}. Furthermore, coupling DW motion with polarized spin-wave excitations offers opportunities in magnonic logic, including spin-wave rotators and interferometers^{31,38}. Finally, geometry-tailored AFM nanostructures fit naturally within the broader framework of curvilinear spintronics, where 3D and bent nanowires are envisioned as building blocks for high-density, low-power information technologies^{33,34,40}.

In conclusion, we have studied the AFM DW dynamics in a helix-shaped geometry driven by a rotating field. We found two motion modes of DW: (i) rigid motion, which is accompanied by the constant DW velocity $\propto \omega$ and fixed phase. The fixing of the DW phase appears due to the curvature-induced DMI term. The DW velocity in the rigid motion is independent of the damping and rotating field strength. The DW velocity for conventional AFM materials like NiO, KFeS₂, and helical single-chain magnets remains on par with spin-orbit torque^{12,49} and coherent spin wave⁵⁰ driven motion. (ii) Oscillating motion appears for the relatively high frequencies ($\omega > \omega_c$) and is accompanied by the drastic drop of DW velocity with ratchet motion. At the same time, the DW phase is oscillating. The value of the critical frequency increases with the curvature increase.

In the limit case of zero curvature, we observe the full synchronization of the field and DW wall rotation frequencies. The direction of the DW phase rotation is defined by the geometrical chirality, i.e., $\Phi = C\omega$. In this limit case, the DW is immobile.

Methods

Collective variable approach

To study the dynamics of the DW we use a collective variable approach based on the generalized $q - \Phi$ model⁵⁸

$$\cos \theta = -p \tanh \frac{\xi - q(t)}{\Delta}, \quad \dot{\phi} = \Phi(t) + a[\xi - q(t)], \quad (8)$$

which is good for the description of DWs in the curvilinear systems^{46,59} and rectilinear systems with intrinsic DMI^{9,58}. Parameters $q(t)$ and $\Phi(t)$ are time-dependent collective variables, which determine the DW position and phase (orientation of the transversal component of Néel vector $\mathbf{n}$ in the center of DW), respectively, $p = \pm 1$ is the topological charge, which determines the DW type: head to head ($p = +1$) or tail to tail ($p = -1$). The DW width Δ and asymmetry of the phase distribution (slope) a are assumed to be slave variables^{9,41,46,59–61}, i.e., $\Delta(t) = \Delta[q(t), \Phi(t)]$ and $a(t) = a[q(t), \Phi(t)]$. One should note that the model (8) corresponds to the exact DW solution in the

helix wires⁴⁶ with equilibrium values $\Delta_0 = 1$ and $a_0 = -\sigma$, for details see (SM) and Fig. 1(c).

Static DW

For the case $\mathbf{h} = 0$, substituting the ansatz (8) in energy density $\mathcal{E}$ (2) and integrating over the arc length, one obtains the total energy of DW

$$\frac{E^{\text{DW}}}{2\sqrt{AK}} = \frac{1+a^2\Delta^2}{\Delta} + \mathcal{K}_1\Delta + \frac{\mathcal{K}_2\Delta}{2} \left[1 + a\pi\Delta \frac{\cos 2\Phi}{\sinh(a\pi\Delta)} \right] + a\mathcal{D}_1\Delta + p\mathcal{D}_2\frac{\pi}{2} (1 + a^2\Delta^2) \frac{\cos \Phi}{\cosh(a\pi\Delta/2)} \quad (9)$$

which is analogous to that for the ferromagnetic thin wire⁴⁶. Minimization of E^{DW} gives

$$\cos \Phi_0 = -\text{sign}(p\kappa), \Delta_0 = 1, a_0 = -\sigma \quad (10)$$

for $\kappa, |\sigma| \ll 1$. The corresponding Néel vector distribution and its schematic representation are shown in Fig. 1(c).

For the case of small curvature and torsion ($\kappa, |\sigma| \ll 1$) the DW energy (9) can be written as

$$\frac{E^{\text{DW}}}{2\sqrt{AK}} \approx \frac{1+a^2\Delta^2}{\Delta} + \Delta + 2\sigma a\Delta + p\pi\kappa \cos \Phi. \quad (11)$$

Limit case of zero curvature

The components of effective Lagrangian and dissipative Rayleigh function (4) can be written as

$$\begin{aligned} \mathcal{L}_0 &= 2\Delta \left(\frac{1+a^2\Delta^2}{\Delta^2} \dot{\Phi}^2 - 2a\dot{\Phi}\Phi + \Phi^2 \right), \\ \mathcal{L}_1 &= -2p\mathcal{C}\pi \frac{\dot{\Phi} + \Delta^2(a+\sigma)(a\dot{\Phi} - \Phi)}{\cosh[\pi\Delta(a+\sigma)/2]} \sin \Psi, \\ \mathcal{L}_2 &= -\pi\Delta^2(a+\sigma) \frac{\cos^2 \Psi}{\sinh[\pi\Delta(a+\sigma)]}, \\ \mathcal{E}^{\text{DW}} &= 2\Delta \left(\frac{1+a^2\Delta^2}{\Delta^2} + 2a\sigma + 1 + \sigma^2 \right), \end{aligned} \quad (12)$$

where we introduce $\Psi = \bar{\omega} - \mathcal{C}\Phi - |\sigma|q$. The Lagrangian and dissipative Rayleigh function (4) with (12) result in equations of motion (5) with the following components of the effective damping matrix $\hat{\mathbb{R}}$ and field-induced force $\mathbf{F}$

$$\begin{aligned} \hat{\mathbb{R}} &= 2p\pi h \cos \Psi \begin{bmatrix} 0 & -1 \\ 1 & 0 \end{bmatrix}, \\ \mathbf{F} &= 2h\mathcal{C} \begin{bmatrix} h\sigma \sin 2\Psi - p\pi\bar{\omega} \cos \Psi \\ h \sin 2\Psi \end{bmatrix}, \end{aligned} \quad (13)$$

where we consider DW phase slope and width with their equilibrium values $a \equiv a_0 = -\sigma$ and $\Delta \equiv \Delta_0 = 1$.

Details of numerical simulations

To verify our analytical calculations we perform a set of numerical simulations for AFM curved wires. We consider a single chain with lattice constant δ_a . Each node is characterized by a magnetic moment $\boldsymbol{\mu}_i(t)$ which is located at the position $\mathbf{r}_i$. Here, $i \in \mathbb{N}$ defines the magnetic moment and its position on the chain with size $i \in [1, N]$. The dynamics of a magnetic system is governed by discrete Landau-Lifshitz-Gilbert equations

$$\frac{d\boldsymbol{\mu}_i}{dt} = \frac{1}{\hbar S} \left[\boldsymbol{\mu}_i \times \frac{\partial \mathcal{H}}{\partial \boldsymbol{\mu}_i} \right] + \eta \left[\boldsymbol{\mu}_i \times \frac{d\boldsymbol{\mu}_i}{dt} \right], \quad i = \overline{1, N}. \quad (14)$$

The Hamiltonian of a magnetic system has the following form,

$$\mathcal{H} = \frac{\mathcal{J}S^2}{2} \sum_{i=1}^{N-1} \boldsymbol{\mu}_i \cdot \boldsymbol{\mu}_{i+1} - \frac{\mathcal{K}S^2}{2} \sum_{i=1}^N (\boldsymbol{\mu}_i \cdot \mathbf{e}_T^i)^2 - g\mu_B S \sum_{i=1}^N \boldsymbol{\mu}_i \cdot \mathbf{H}, \quad (15)$$

where $\mathcal{J} > 0, \mathcal{K} > 0$ are the exchange and anisotropy constants, respectively, $\mathbf{H}$ is the external magnetic field which rotates in xy plane with frequency (ω , and $\mathbf{e}_T^i$) is a tangent unit vector to the wire. Note that the vector $\mathbf{e}_T^i$ introduces the information about the wire shape into the model (15).

The magnetization dynamics is simulated by means of numerical solution of the set of ordinary differential equations (14) for the initial conditions determined by the initial magnetization.

Data availability

Data is provided within the manuscript or supplementary information files.

Code availability

Numerical code available upon reasonable request.

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Author contributions

D.D.S. supervised the project. K.V.Y. and D.D.S. formulated the theoretical problem. Y.A.B. performed spin-lattice simulations. Y.V.B. and K.V.Y. performed analytical calculations. The manuscript was written by K.V.Y. with contribution from D.D.S. and Y.A.B.

Competing interests

The authors declare no competing interests.

Additional information

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Correspondence and requests for materials should be addressed to Denis D. Sheka.

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