Intrinsic chiral field as vector potential of the magnetic current in the zig-zag lattice of magnetic dipoles

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Chiral magnetic insulators manifest novel phases of matter where the sense of rotation of the magnetization is associated with exotic transport phenomena. Effective control of such phases and their dynamical evolution points to the search and study of chiral fields like the Dzyaloshinskii–Moriya interaction. Here we combine experiments, numerics, and theory to study a zig-zag dipolar lattice as a model of an interface between magnetic in-plane layers with a perpendicular magnetization. The zig-zag lattice comprises two parallel sublattices of dipoles with perpendicular easy plane of rotation. The dipolar energy of the system is exactly separable into a sum of symmetric and antisymmetric long-range exchange interactions between dipoles, where the antisymmetric coupling generates a nonlocal Dzyaloshinskii–Moriya field which stabilizes winding textures with the form of chiral solitons. The Dzyaloshinskii–Moriya interaction acts as a vector potential or gauge field of the magnetic current and gives rise to emergent magnetic and electric fields that allow the manifestation of the magnetoelectric effect in the system.

Chiral symmetry refers to symmetry under mirror reflection: a subject is said to be chiral when it lacks such symmetry. Chiral asymmetry is rather common in nature at several scales: at the microscale, it is well-known that elementary particles and organic molecules have a preferential chiral characterization, while very large macroscopic systems like spiral galaxies are chiral too¹. In condensed matter physics, a theory, symmetry, or field is chiral if it is not invariant under the inversion of one spatial dimension. Chiral condensed matter systems realize exotic electronic², topological³ and magnetic phases⁴–¹¹ and structures¹²,¹³. The transport properties of chiral matter can undergo quantum anomalies associated with chiral symmetry as in the topological Dirac/Weyl semi-metals¹⁴, the antisymmetric Dzyaloshinskii–Moriya interaction¹⁶, (DMI), is a chiral coupling¹⁷ able to trigger magnetic torques in magnetic systems. Such torques can stabilize localized and spatially modulated structures with a fixed sense of rotation¹⁸ and influence the transport in such systems; in consequence, they are key to the development of spin-based memory, logic, and signal transmission devices¹¹. DMI is generally described by a vector product formed by the localized magnetic moments of two magnetic ions \( U_{DM} = D_{ij} \cdot (\mathbf{m}_i \times \mathbf{m}_j) \), where \( D_{ij} \) is called Dzyaloshinskii–Moriya (DM) vector. Dzyaloshinskii¹⁹ introduced \( U_{DM} \) based on phenomenological considerations to explain the observation of weak ferromagnetism in some antiferromagnets. Later, Moriya²⁰ demonstrated that in low symmetry magnetic crystals, the spin-orbit interaction can lead to DMI by taking into account the spin-orbit coupling (SOC) in the theory of superexchange interaction proposed by Anderson²¹. In this theory \( D_{ij} \) is proportional to the spin-orbit interaction and depends on the symmetry of crystals. Afterward, it was shown that this mechanism is relevant only when the local symmetry is sufficiently low and that a weak ferromagnetic moment emerges from the superexchange coupling only when more than a single bond is considered²². Chiral magnetic couplings have been induced in centrosymmetric crystals by symmetry breaking due to electric currents²³, applied magnetic and electric fields and by strain²⁴. For bulk magnetic materials, such coupling is generally weak; however, in small artificial structures such as ferromagnetic films, multilayers, nanowires, and nanodots, this is not the case²⁵–³². In two-dimensional films, the

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interfacial DMI defines a rotational sense for the magnetization which can be used to create chiral magnetic structures like spin-spirals, domain walls, and skyrmions\textsuperscript{33,34}. Recently, a strong interlayer Dzyaloshinskii–Moriya interaction has been demonstrated in FM/Pt/FM trilayers with orthogonal magnetization. In this system, the DMI causes a chiral interlayer coupling\textsuperscript{35–37} that favors one-handed orthogonal magnetic configurations of Co and TbFe, as revealed through the Hall effect and magnetoresistance measurements\textsuperscript{38}. At the nanoscale, the intralayer DMI has been used to engineer strong, localized intrinsic chiral torques that trigger the spontaneous motion of domain walls or bias the speed of current-driven domain walls in the magnetic race-track\textsuperscript{39–42,44}. In thin-film metallic systems, spin-orbit coupling arises from a proximal heavy metal\textsuperscript{43}, where the metallic layer induces an electric field orthogonal to the plane of the lattice\textsuperscript{38,40,41}. Moreover, recent findings in perpendicularly-magnetized iron garnets, rare-earth orbital magnetism has given rise to an intrinsic spin-orbit coupling generating interfacial DMI at mirror symmetry-breaking interfaces\textsuperscript{45}. Moreover, recent findings showing that the rare-earth ion substitution and strain engineering can significantly alter the DMI\textsuperscript{44,46,47}, remain to be understood.

**Summary of results.** Aimed to identify new sources of chiral fields and stable chiral structures, we study the magnetization dynamics of macroscopic zig-zag lattices of dipoles with experiments and theory. Magnetic dipoles in different sublattices have perpendicular easy planes of rotation, which, combined with dipolar interactions, allows the exact mapping of the magnetic energy into four energetic contributions, which include symmetric and antisymmetric or chiral long-range interactions between the dipoles. Dynamics in the system is induced by tuning the chiral torques through the variation of the gap $\ell$, which is the distance between the two sublattices along axis $\hat{y}$ (Fig. 1). As $\ell$ is varied, the system transits between four magnetic phases through a rich dynamical process that features hysteresis and stabilizes chiral magnetic textures. The explicit formula for symmetric and chiral couplings exposes the DMI vector acts as a vector potential of an out-of-plane magnetic current. Such gauge potential gives rise to a magnetic field and a magnetic flux which, through variations of $\ell$, induces an electric field orthogonal to the plane of the lattice\textsuperscript{38,40,41}.

The paper is organized as follows. In section “Model” we present the model system and show the separation of the total energy of it into symmetric and antisymmetric contributions. Section “Magnetic phases in terms of $\ell$” shows the magnetic phases realized by the lattice as $\ell$ is tuned. In section “Symmetric and antisymmetric contributions to the dipolar energy” we discuss the contributions of the effective symmetric and antisymmetric couplings to the energetics of the system, and in section “Intrinsic magnetic Current and torque” we define the magnetic current and torque in terms of the chiral field that arises product of the dipolar interactions. Section “Magnetic current in the planar states, $\ell_f < \ell < \ell_c^*$” is devoted to studying the magnetic current, the associated potential vector, and the emergent fields that arise in the planar phase. Section “Chiral soliton at $\ell > \ell_c^*$” focuses on the magnetic phase realized at large $\ell$ and the onset of magnetic solitons. Concluding remarks are presented in section “Conclusions”.

**Model.**

The magnetic dipolar energy for the system of $n$ dipoles in the zig-zag lattice reads $U = \frac{g}{2} \sum_{i \neq j=1}^{n} \left( \frac{\mu_0 m_i \cdot m_j}{|r_i - r_j|^3} \right)$, where $\hat{m}_i = (r_i - r_j)/|r_i - r_j|$ and $g = \frac{\mu_0 m_0^2}{4\pi\Delta}$ sets the energy scale. It contains the physical parameters involved in the energy, such as $\Delta$, the lattice constant, $\mu_0$, the magnetic permeability, and $m_0$, the intensity of the magnetic moments with saturation magnetization $M_s$. Hereafter we normalize all distances by $\Delta$. The magnetic moments are normalized by $m_0$ and dipoles belonging to sublattice $\alpha$ : $(c, p)$ have unit vector $\hat{m}_i^{\alpha} = (\sin \theta_i^{\alpha} \cos \phi_i^{\alpha}, \sin \theta_i^{\alpha} \sin \phi_i^{\alpha}, \cos \theta_i^{\alpha})$. The $n$ dipoles are located at the vertices of a zig-zag lattice made out of two sublattices that are coplanar parallel chains: chain $c$ with $n_c$ dipoles and chain $p$ with $n_p = n_c - 1$ as shown in Fig. 1. Dipoles rotate in an easy plane in terms of a polar angle $\theta$ with respect to the $\hat{z}$ axis, and a fixed azimuthal angle $\phi = 0$ and $\phi = \frac{\pi}{2}$. Hence, dipoles in $c$ rotate in the $\hat{x} - \hat{z}$ plane and dipoles in $p$ in the $\hat{y} - \hat{z}$ plane (Fig. 1). With easy planes mutually perpendicular among chains, the dipolar energy is exactly separable into symmetric and antisymmetric long-range interactions:

$$U = \frac{g}{2} \sum_{i \neq j=1}^{n} \left[ J_{ik}^S \left( \frac{3}{2} \cos(\theta_i^c + \theta_j^c) - 1 \right) + \hat{m}_i^c \cdot \hat{m}_j^c + J_{ik}^A \left( \hat{m}_i^p \cdot \hat{m}_j^p + \hat{m}_i^c \cdot \hat{m}_j^c + \hat{\sigma}_{ik} \cdot (\hat{m}_i^c \times \hat{m}_j^c) \right) \right]$$

which gives rise to four energetic contributions to the full magnetic energy of the system, consecutively denoted such that $U = U_{S} + U_{p} + U_{AP} + U_{DM}$. They correspond respectively to symmetric intra-sublattice interactions in $c$ and $p$, a symmetric inter-sublattice interaction, and an antisymmetric inter-sublattice interaction energy. Explicit formulas for the associated couplings read, $J_{ik}^S = \frac{1}{|r_i - r_j|}$, $J_{ik}^A = \frac{1}{(i^2 + (i-k+\frac{1}{2})^2)^{3/2}}$, which are respectively symmetric intra-chain and interchain couplings. $\hat{\sigma}_{ik} = -3 \left( \begin{array}{c} 0, 0, 0 \\ 0, -\frac{\ell(i-k+\frac{1}{2})}{(i-k+\frac{1}{2})^2 + \ell^2} \end{array} \right)$ corresponds to an interchain Dzyaloshinskii–Moriya antisymmetric coupling, perpendicular to the plane of the lattice. Note the dependence of $J_{ik}^A$ and $\hat{\sigma}_{ik}$ on $\ell$. 

https://doi.org/10.1038/s41598-023-28545-9
Figure 1. (a) Magnetic phases of the experimental zig-zag lattice in terms of \( \ell \). From top to bottom as \( \ell \) decreases: at large \( \ell > \ell_* \), the system settles in the Tw phase. In the range \( (\ell_* , \ell) \) the lattice enters the metastable regime in the \( x-y \) plane AF2/FAF. As \( \ell \) decreases further, in the range \( \ell \in (\ell_f , \ell) \), the AF2 phase is selected. At very small gaps \( \ell < \ell_f \) dipoles of both sublattices configure an antiferromagnetic state along the \( \hat{z} \) axis. Inset: The zig-zag lattice's geometry consists of \( n = 37 \) Neodymium magnets hinged on top of a PTFE plate. All have length \( a = 12.7 \times 10^{-3} \) [m], radius \( r = 0.79 \times 10^{-3} \) [m], mass \( 0.189 \times 10^{-3} \) [kg], and saturation magnetization \( M_s = 1.05 \times 10^6 \) [A/m]. The distance between two consecutive rods in the same chain is \( \Delta = 22 \times 10^{-3} \) [m] fixed, and the tunable vertical interchain gap is \( \ell = y/\Delta \) where \( y \) is the vertical distance measured from chain \( c \). Dipoles at chain \( c \) and \( p \) rotate in the planes \( x-z \) and \( y-z \) respectively. (b) Screenshots of the lattice from Molecular dynamics simulations (see supplemental material 1 for details).
Magnetic phases in terms of ℓ

The experimental setup comprises $n = 37$ Neodymium cylindrical magnets of length $a$, radius $r$ and $m_0 = \alpha r^2 \mu_0 M$, hinged at the sites of a Polytetrafluoroethylene (PTFE) plate forming a zig-zag lattice with lattice constant $\Delta$. Sublattices $c$ and $p$ have respectively $n_c = 19$ and $n_p = 18$ magnets (Fig. 1) and rotate in the mutually perpendicular planes $x - z$ and $y - z$ respectively. A small amount of disorder due to deviations of the dipoles with respect to their easy plane of rotation follows a Gaussian distribution centered at zero and with a standard deviation $\delta \phi \sim 0.005$. The interchain gap $\ell$ is tuned by moving chain $p$ along the $y$ axis in the range $\ell \in (0, 1.5)$ at a constant speed $v$ while $c$ remains at rest. A camera captures magnetic configurations of the system as the stage with chain $p$ is moved from $\ell = 1.5$ to 0 (approaching) and back to $\ell = 1.5$ (receding). In Fig. 1a we indicate the north poles of the magnets with a black tip.

Depending on $\ell$, dipoles settle into four magnetic configurations as shown in Fig. 1a. At small gaps $\ell \in (0, 0.2) = (0, \ell_f)$ the system realizes an out-of-plane antiferromagnetic parallel phase, AF, along $\hat{z}$. It consists of dipoles arranged ferromagnetically with respect to the others in the same sublattice and antiferromagnetically with respect to dipoles in the other sublattice. Increasing $\ell$ triggers a spin flop transition where all dipoles depart from $\hat{z}$ and settle in the $x - y$ plane featuring phase AF$^2$ ($\ell \in (0.2, 0.6) = (\ell_f, \ell)$), where chain $c$ describes a collinear antiferromagnetic state, and chain $p$ a parallel antiferromagnetic configuration. For intermediate interchain gaps, $\ell \in (0.6, 0.8) = (\ell, \ell_c)$, phase AF$^2$ competes with phase FAF which differs from AF in that chain $c$ arranges in a ferromagnetic collinear fashion. At large gaps, $\ell_c < \ell < \ell < 1$ chain $c$ settles in a collinear ferromagnetic state, while chain $p$ defines a winding texture consisting of a twisted parallel antiferromagnetic state in the $y - z$ plane. This phase is denoted Tw. The previous magnetic configurations define the magnetization curves shown in Fig. 2. At small gap $\ell \in (0, \ell_f)$, Fig. 2a shows the absence of average magnetization along $\hat{x}$, $M_x$. Instead, in each sublattice the magnetization along $\hat{z}$ has a different sign and reaches its maximum in this regime (bottom panel of Fig. 2b and top panel of Fig. 2c). Finite staggered magnetizations along $\hat{x}$, $N_x$ (Fig. 2b top panel) for this size of the gap shows that dipoles in $c$ realize a canting state in the $x - z$ plane, consistent with Fig. 1. Indeed, the AF magnetic state is difficult to achieve in experiments due to the strong dipolar interactions among nearest neighbor dipoles in different sublattices, the azimuthal disorder and the frictional rotation of the magnets. For $\ell \in (\ell_f, \ell_c)$ dipoles are in the $x - y$ plane so that $M_x$ goes to zero in both sublattices and staggered magnetizations $N_x$ (not shown) reach their maximum values, Fig. 2b. The metastable regime (see supplemental material 1 for details) with competing phases AF$^2$ and FAF originates the hysteresis loops of Fig. 2 at intermediate $\ell \in (\ell, \ell_c)$. Finally, at large values of the gap $\ell_c < \ell < 1$, chain $c$ remains ferromagnetic which results in $M_x = 1$ and the winding texture in $p$ is such that $N_x = 0$ and $N_x^c \neq 0$ while $M_y = 0$. We further examined the evolution of the system with $\ell$ by implementing molecular dynamics simulations (see supplemental material 1 for details). Magnetic phases from numerics coincide with those found in experiments as shown by the screenshots of the numerical lattice in Fig. 1b. The hysteretic behavior seen in experiments at intermediate $\ell$ is captured too by simulations as can be verified in the magnetization loops of $M_x$ and $N_x$ shown in supplementary Fig. 1. The width of the loops is well reproduced by considering a nearest neighbor interacting model (see supplemental material 1 for details). In addition to the loop at $\ell \in (\ell_f, \ell_c)$, numerics reveals another loop for $\ell \in (\ell_c, \ell_f)$. It shows that phases AF$^2$ and AF are metastable in this regime. This is consistent with the spin flop transition being of the first order type$^{47,48}$. In the limit of a large gap, $\ell > 1$ we find that the twisted state relaxes into the FAF phase.

Figure 2. Magnetization dynamics in experiments. (a) evolution of $M_x$ and (b) (top panel) the staggered magnetization along $\hat{x}$, $N_x$ in terms of $\ell$. Red arrows pointing to the right and left denote chain $p$ moving apart from and approaching chain $c$ respectively. The bottom panel of (b) and the upper panel of (c) show the $\hat{z}$ component of the magnetization of the $p$ and $c$ sublattices, respectively. The magnetic current along $\hat{x}$ in the metastable regime at intermediate $\ell$ is shown in the bottom panel of figure (c).
Symmetric and antisymmetric contributions to the dipolar energy

Equation 1 reveals the specific contribution of the symmetric and antisymmetric inter and intra-sublattice interactions to the system’s total energy. Figure 3 shows the evolution of each of them with \( \ell \) in units of \( \frac{g^2}{2} \). In Fig. 3a (experiment) and c (numerics) the minimum of the total dipolar energy \( U \) occurs at the onset of phase AF. This extreme in \( U \) coincides with the optimum of \( U_{\text{DM}} \) (blue curves in Fig. 3b,d) and with the maximum of the spin current along the \( \hat{z} \) direction as shown in the bottom panel of Fig. 2c. Being proportional to the \( \hat{z} \) projection of dipoles in chain \( c \), the energy \( U_{\text{cp}} \) (in black) is non null during phase AF and it is the dominant contribution to \( U \) at \( \ell < \ell_f \). \( U_{\text{DM}} \) (in blue) dominates the dynamics followed by the intrachain symmetric energy \( U_c \) (red) which becomes the dominant contribution to the energy once the system is in the Tw phase. Finally, Fig. 3c,f show the intra-chain symmetric energy in \( p \), \( U_p \) to be non null and one order of magnitude smaller than the others in all the range of \( \ell \). At \( \ell > \ell_s \) the system is in the Tw phase and the total energy, dominated by \( U_c \), barely changes with \( \ell \). Note that in this magnetic phase the dipoles arrange such that the interchain energy contributions cancel out \( U_{\text{DM}} = U_{\text{cp}} = 0 \). Experiments (Fig. 3a–c) and numerics (Fig. 3d–f) agree fairly well. Next, consider the case of a long zig-zag lattice. We denote the horizontal distance between two dipoles \( x \) and the vertical gap \( y \). The Dzyaloshinskii–Moriya coupling perpendicular to the plane of the system is written as,

\[
D = -\frac{3}{2} \frac{y(x+\frac{1}{2})}{(x+\frac{1}{2})^2+y^2} \hat{z}
\]

and decays fast with dipole distance (see supplementary Fig. 6 (see supplemental material for details)). \( D \) reaches a maximum for nearest neighbor dipoles \( (x = 0) \) at an optimum interchain distance \( y_m = \frac{1}{4}(1 + 2x) = \frac{1}{2} \) and its contribution to the total energy is comparable with that of the symmetric energies (see supplemental material 1 for details). Integrating out the \( x \) coordinate yields an explicit formula for the effective gap-dependent interchain chiral coupling in the system,

\[
D_{\text{ef}} = \frac{8}{3} \frac{y}{\left(4y^2+1\right)^{3/2}} \quad (y>0).
\]

Similar to the previous case, the symmetric inter and intrachain couplings, \( J \) and \( J^0 \) decay very fast with \( x \) and \( y \), but at small \( (x, y) \), \( D \) becomes the largest among the three. The formula for the effective gap dependent interchain symmetric coupling after integrating out the \( x \) coordinate yields

\[
J_{\text{ef}} = \frac{1 - \sqrt{4y^2+1}}{y} \quad (y>0).
\]

As shown in supplementary Fig. 7 (see supplemental material for details), \( J(\ell)_{\text{ef}} > D(\ell)_{\text{ef}} \) for \( \ell < \ell_{nf} \) point at which they are equal and from there \( D(\ell)_{\text{ef}} \) becomes the leading coupling.

Figure 3. (a,d) show the total energy of the system as a function of \( \ell \) in experiments and numerics, respectively. The magnetic phases are delimited by the dotted vertical lines. Using Eq. 1 the contribution of symmetric and antisymmetric long-range interactions to the total magnetic energy is shown in (b) (experiments) and (e) (numerics) with \( U_c, U_{\text{cp}} \) and \( U_{\text{DM}} \) shown in red, black and blue respectively. The dark blue curves of (c) and (f) show \( U_p \) in experiments and numerics, respectively.
Intrinsic magnetic current and torque

The spin current arises from the Heisenberg equation\(^49\) whose correspondence to the classical system at hand is:
\[
\frac{\partial \mathbf{m}}{\partial t} = \mathbf{m} \times \frac{\partial \mathbf{H}}{\partial \mathbf{m}} = \mathbf{m} \times \mathbf{H}_i = -\sum_{\ell} \mathbf{f}_{ik},
\]
where \(\mathbf{H}_i\) denotes the internal magnetic field produced by all dipoles but the \(i\)-th at the position of \(\mathbf{m}_i\), and \(\mathbf{f}_{ik}\) is the associated torque. Hence, a magnetic current is induced by the internal magnetic torque, which can be tuned by changing \(\ell\). Writing down \(\mathbf{U}\) in terms of \(\mathbf{f}_{ik}\), the interaction matrix of the system, (see supplemental material 1 for details) yields \(\mathbf{U} = -\frac{1}{2} \sum_{i,k} m_{ik} \mathbf{f}_{ik} m_k\), which allows to formulate the classical correspondence of the spin current in terms of \(\mathbf{f}_{ik}\) as follows,
\[
\mathbf{J}_{ik}^{(a)} = \epsilon^{abc} m_i^{(b)} \mathbf{f}_{ik}^{(c)} m_k^{(d)}. \tag{2}
\]
In Eq. (2) \(a\) labels the vector components \(x, y, z\) and the matrix elements of \(\mathbf{f}_{ik}\) correspond to the magnetic interactions that couple dipoles in Eq. (1) as shown in supplemental material 1. \(\epsilon^{abc}\) denotes the Levi-Civita symbol, and thus Eq. (2) demonstrates that a magnetic current arises from the matrix elements of \(\mathbf{f}_{ik}\) connecting magnetization vectors of dipoles coupled by the chiral DM vector \(\mathbf{D}\).

Magnetic current in the planar states, \(\ell_f < \ell < \ell_\star\).

Figure 2 shows the \(z\) component of the magnetic current, \(\mathbf{J}^{(z)} = 2 \sum_{i,k,a} \sum_{k=0}^\infty \mathbf{f}_{ik} m_i^m m_k^m\). \(\mathbf{J}^{(z)}\) connects magnets whose magnetization have perpendicular components in the \(\hat{x} - \hat{y}\) plane. It is zero or negligible at \(\ell < \ell_f\), and after reaching its maximum in phase AF\(^2\) it becomes zero once again in phase Tw. As expected from the magnetization loops, \(\mathbf{J}^{(z)}\) realizes hysteresis in the metastable regime (\(\ell < \ell_\star\)).

\(\mathbf{D}\) as a vector potential to the magnetic current.

Using Eq. (2), the magnetic current in phase AF\(^2\) yields,
\[
\mathbf{J}_{AF^2} = 2g \left[ 0, 0, \sum_{i,k,a \neq \beta} \mathbf{D}_{ik} \sin(\theta_a^i - \theta_a^k + \frac{\pi}{2}) \right]. \tag{3}
\]
where \(2g\mathbf{D}_{ik}\) is interpreted as the spin stiffness or magnetic rigidity of the system\(^49\). In phase AF\(^2\) the energy of the system \(U_{AF^2} = U_c + U_p + U_{DM}\) (Eq. 1) can be written,
\[
U_{AF^2} = \sum_{i,k,\alpha \neq \beta} \left[ 8\alpha^p f_{ik}^{\alpha^p} + (1 - 8\alpha^p) \mathbf{D}_{ik} \frac{\pi}{2} \right] \sin \theta_a^i \sin \theta_a^k \tag{4}
\]
where \(f_{ik}^{\alpha^p} = -2f_{ik}^{\alpha^p}\) and \(f_{ik}^{\beta^p} = f_{ik}^{\beta^p}\). It can be written as
\[
u_{AF^2} = \sum (x_{ik}^{\alpha^p} \cos A_{ik}) \sin \theta_a^i \sin \theta_a^k \tag{5}
\]
where \(x_{ik}^{\alpha^p} = \cos A_{ik} = \left(8\alpha^p f_{ik}^{\alpha^p} + (1 - 8\alpha^p) \frac{\mathbf{D}_{ik}}{\pi} \right)\) (see supplemental material 1 for details). Correspondingly, the magnetic current becomes
\[
\mathbf{J}_{AF^2} = 2 \sum x_{ik}^{\alpha^p} \sin \left( \theta_a^i - \theta_a^k + \frac{\pi}{2} + A_{ik} \right) \tag{6}
\]
Consequently, the chiral DM vector \(\mathbf{D}\) acts as the vector potential or gauge field associated to the magnetic current.\(^49\). Further, \(\mathbf{D}\) gives rise to the magnetic field along the \(\hat{x}\) axis, \(\mathbf{H}_{DM} = \sum_i \mathbf{m}_i \times \mathbf{D}_{ik}\) which acts on dipole \(i\) in the \(c\) sublattice. In a large zig-zag lattice \(\mathbf{H}_{DM}\) acts as an effective interchain field \(\mathbf{H}_{DM} = N_y \times \mathbf{D}\) whose magnitude is proportional to the staggered magnetization along \(\hat{y}\) and to the chiral DM coupling. This effective field \(\mathbf{H}_{DM}\) originates a magnetic flux \(\Phi_{DM}\) in a loop \(S\) parallel to the \(y - z\) plane. Because \(P\) moves at a speed \(v = \frac{dy}{dt}\) with respect to \(c\), the time derivative of \(\Phi_{DM}\) induces a femtovolt \(E_{DM} = 2\pi v\bar{y}H_{DM}\) and an electric field \(E_{DM}\) in such a loop due to Faraday’s law. The induced electric field points along the \(\hat{z}\) direction and is proportional to the relative speed between sublattices and the magnitude of the DM vector, \(E_{DM} \sim v\mathbf{D}_N = vE_{DM}\) (see supplemental material 1 for details). This leads to the coupling between the magnetic current and the induced electric field through the electric polarization. \(P = \frac{E_{DM}}{\bar{y}}\). The control of electric polarization by magnetic fields in a magnetic system is known as the magnetoelectric effect.\(^50\). Here we show a possible route for the manifestation of magnetoelectricity at magnetic interfaces. Furthermore, the existence of an intrinsic electric field in the purely magnetic zig-zag lattice could be important for two reasons: 1) it can be manipulated by changing the DM coupling (proportional to \(g/m_0\)), which is tuned by changing the geometric parameter \(\ell\), and 2) it has been predicted that electric fields could be able to propel magnetic topological textures in antiferromagnets.\(^51\).

Chiral soliton at \(\ell > \ell_\star\).

The nonlinearity of the spin dynamics of magnets is primarily determined by the purely geometric properties of the magnetization field that one sublattice exerts on the other. These properties give rise to topologically non-trivial structures in the zig-zag chain of dipoles. At \(\ell_\star < \ell < 1\), dipoles belonging to different sublattices remain orthogonal with respect to each other, and thus the equilibrium orientation of the dipoles minimizes \(U_{AF^2}\). Correspondingly, in this regime, the magnetic current \(\propto \sin (\theta_a^i - \theta_a^k + \frac{\pi}{2} + A_{ik})\) remains zero. Once sublattice \(c\) settles into the ferromagnetic collinear state at \(\ell_\star\), the DM energy can be rewritten via the internal effective Dzyaloshinskii field, felt by dipoles in \(p\) due to their antisymmetric interaction with collinear ferromagnetic dipoles in \(c\): \(\mathbf{D} (x, \ell, M_c\bar{y})\). When the two sublattices are not farther than \(\Delta\), this field destroys the parallel
antiferromagnetic state of \( p \) because the Zeeman energy orient dipoles along the Dzyaloshinskii field produced by \( c \). Further, the effective Dzyaloshinskii field acts as an anisotropy internal field that rotates dipoles in \( p \) along the \( \hat{x} \) axis, giving rise to a transverse magnetization along \( \hat{z} \). Consider the Neel order in \( p \) constrained in the \( y-z \) plane, \( N_p = (0, \sin \theta, \cos \theta) \). The angle \( \theta(x,t) \) parametrizes the local magnetic state. As shown in Fig. 3, in phase \( Tw \) the full energy remains constant while \( U_{DM} = U_{pp} = 0 \). Therefore \( \theta(x,t) \) is such that minimizes \( U_c + U_p \). Because of the fast decay of the couplings \( D \) and \( J \) (see supplemental material 1 for details) we consider interchain interactions up to second nearest dipoles. Denoting \( \omega_1 = (J(1,1) + D(1,1)) \) and \( \omega_2 = (J(2,1) + D(2,1) + \sum_{j=1}^\infty f^0(x)) \) (see supplemental material 1 for details), the magnetic \( m_p \) and \( m_q \) at each position of \( p \) are shown in Fig. 4. The solution to the equation \( \frac{\partial \theta}{\partial t} = 2 \frac{\omega_1}{\omega_2} \) which yields \( \theta(x,t) = 2 \arctan \left( \frac{\omega_1}{\omega_2} \right) \) that corresponds to a one-dimensional soliton. Evolution of the twisted structure as \( \ell \) grows from \( \ell_\ast \) up to 1 and the values of \( m_p \) and \( m_q \) at each position of \( p \) are shown in Fig. 4. The soliton consists of two Bloch domain walls (each domain wall is realized by one or more dipoles that have rotated toward the \( \hat{z} \) axis) which are born near the edges of \( p \) at \( \ell \sim \ell_\ast \). The net total topological charge is a conserved quantity, and the associated continuity equation defines the dynamics of the winding texture as \( \ell \) is tuned, as shown in Fig. 4. As \( \ell \) grows from \( \ell_\ast \), the two domain walls extend toward the center of \( p \), including more dipoles, until they merge. Once they merge at \( \ell \sim 1 \), the domain walls disappear, and the sublattice is such that all the dipoles orient along the same axis. These winding structures are the product of the internal chiral field and have associated a handedness which determined by the sense of \( M_c \) in \( c \).

Conclusions

We have shown that the dynamics of a zig-zag lattice of dipoles is induced by a magnetic torque, which arises only due to an intrinsic DMI between its sublattices without the aid of external sources to break time-reversal symmetry. The hysteretic dynamics in the system is propelled by interlayer gap variations that tune the internal chiral and achiral fields. The Dzyaloshinskii–Moriya interaction acts as the vector potential of the magnetic current perpendicular to the plane of the lattice, inducing magnetic and electric fields, which allows the manifestation of the magnetoelectric effect at magnetic interfaces. We identified a gauge or vector potential in the system which depends on the interchain couplings. The curl of such vector potential gives rise to magnetic fields and the time derivative of these internal magnetic fields induces electric fields. These intrinsic fields satisfy the Maxwell equations. Therefore in this sense, they are analogs to E and B of the usual U(1) gauge theory of electromagnetism. Such emergent fields are tuned by changing \( \ell \) the distance between the two sublattices. Atomic-scale or mesoscopic spin textures with all broken mirror symmetries and preserved time-reversal symmetry like the twisted magnetic order shown here provide a promising platform to study cross-coupled ferroic orders, magnetic optical activities, and topological transport properties.

Methods

The experimental setup consists of one fixed stage and a computer-controlled movable stage, machined from acrylic plexiglass and covered with Teflon to reduce rotational friction. A camera (Nikon D750) records the rotation of the magnets that are free to move in their respective planes due to carbon fixtures. The NdFeB magnets

\[
\omega(x) = \frac{\omega_1}{\omega_2} \frac{\partial \theta}{\partial t} = 2 \arctan \left( \frac{\omega_1}{\omega_2} \right)
\]

Figure 4. Solitons in the \( p \) sublattice. From top to bottom the left panel shows the winding texture for different values of \( \ell \) as it grows from \( \ell \sim \ell_\ast \) up to \( \ell \sim 0.95 \). The right panel shows the magnetization of each dipole in \( p \) along \( y \) and \( z \) directions.
(grade N42) have a Nickel coating, a radius \( r = 0.79 \pm 0.01 \times 10^{-3} \text{[m]} \), length \( a = 12.7 \pm 0.01 \times 10^{-3} \text{[m]} \), and mass \( 0.189 \times 10^{-3} \text{[kg]} \). The speed of the translation stage was set to 0.1 \text{[mm/s]} \), while the camera records 59 frames per second. These images are converted to rotation angles using standard imaging techniques.

To compute the damping coefficient \( \eta \), we isolated a single rotor, impulsively applied torque to it, and then recorded its relaxation dynamics. We extracted the evolution of \( \theta(t) \) using standard imaging techniques corresponding to damped dynamics without external forcing. The damping is computed directly by fitting it to the solution \( \theta = \exp \left( -t/\tau_D \right) \), and thus we estimated the damping time of a single rod to be \( \tau_D = 0.83 \pm 0.18 \text{[s]} \).

**Data availability**

The datasets used and analyzed during the current study are available from the corresponding author upon reasonable request.

Received: 11 November 2022; Accepted: 19 January 2023

Published online: 23 January 2023

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**Acknowledgements**

P.M. acknowledges support from Fondecyt under Grant No. 1210083. K.H. acknowledges support from the Leading House for the Latin American Region (CLS-HSG) Mobility Grant 2020 and the Swiss National Science Foundation (Projects No. 200020 172774).

**Author contributions**

P.M. conceived the experiment, performed calculations, numerical simulations and wrote the manuscript. A.C. and K.H. conducted the experiments, I.T. performed simulations. All authors analysed the results and reviewed the manuscript.

**Competing interests**

The authors declare no competing interests.

**Additional information**

**Supplementary Information** The online version contains supplementary material available at [https://doi.org/10.1038/s41598-023-28545-9](https://doi.org/10.1038/s41598-023-28545-9).

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