Phase-locking of magnetic vortices mediated by antivortices

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Synchronized spin-valve oscillators may lead to nanosized microwave generators that do not require discrete elements such as capacitors or inductors. Uniformly magnetized oscillators have been synchronized, but offer low power. Gyrating magnetic vortices offer greater power, but vortex synchronization has yet to be demonstrated. Here we find that vortices can interact with each other through the mediation of antivortices, leading to synchronization when they are closely spaced. The synchronization does not require a magnetic field, making the system attractive for electronic device integration. Also, because each vortex is a topological soliton, this work presents a model experimental system for the study of interacting solitons.

A magnetic vortex is a flux-closure magnetic state that consists of a curling in-plane magnetization and an out-of-plane vortex core. It can represent the lowest-energy configuration of a single-domain magnet or be nucleated by a circular magnetic field. This magnetic configuration shows a very rich dynamic behaviour. Among the accessible modes that may be excited, the circular motion of the core around its equilibrium position, called gyrotropic precession, has attracted a large amount of interest. This mode can be excited by a spin-polarized current through the transfer of spin-angular momentum to the local magnetic moment, a phenomenon called spin transfer torque.

Spin transfer torque has been routinely used to control the orientation of magnetic nanostructures. It can also drive the magnetic orientation of a layer into an oscillation, creating a spin transfer nano-oscillator (STNO). Its use in exciting a magnetic vortex in spin-valve nanocontacts has attracted interest because in this vortex-based STNO, a large part of the magnetoresistance is transformed into power. As a consequence, the power is much larger than that generated by uniform-mode STNOs, although the maximum frequency is lower (<2 GHz). The synchronization of two closely spaced STNOs with uniform excitation modes has been demonstrated through spin-wave coupling, leading to significant improvements in their output features.

Here we demonstrate that local synchronization is possible between vortex-based STNOs. We show that vortices orbiting beneath closely spaced nanocontacts can synchronize through the mediation of antivortices. As in the case of phase-locking between uniform modes, the phase-locked state is characterized by decreased linewidth and increased power. Remarkably, the best condition for phase-locking in the system presented here is the absence of externally applied magnetic field. The mechanism of phase coupling is interpreted by using a simple analytical model and on the basis of the results of micromagnetic simulations for the dynamics of interacting solitons and antivortices.

An integrated array of nano-oscillators

In our study we nucleated magnetic vortices under a 2 × 2 array of nanocontacts with centre-to-centre spacing d, as shown in Fig. 1. We prepared devices with d = 10 μm, 2 μm and 500 nm. The magnetic structure was a sputtered deposited multilayer consisting of a 5-nm tantalum/40-nm copper/4-nm permalloy (Py) bottom electrode, a 6-nm gold spacer and a 15-nm cobalt/100-nm gold top electrode. The thickness of the ferromagnetic layers were chosen in such a way as to favor nucleation of vortices in the cobalt layer. The nanocontacts were opened by conductive-tip atomic force microscopy (AFM) nano-indentation and plasma etching into a 40-nm resist spun on the spacer layer before depositing the top electrode. A diameter of ~20 nm could be estimated for the contacts by AFM scan. The lithographically defined top electrode connected the four contacts in parallel. The dynamics of this system could be probed by measuring the induced radio-frequency signal produced at the electrodes because of the giant magnetoresistance effect.

In Fig. 2a, we display the power spectral density (PSD) for the device with intercontact distance d = 2 μm in an H_{bias} = +50 mT perpendicular applied field. The field sign is defined according to the polarity of a previously applied large reset field (see Methods). To stabilize the initial magnetic configuration with a vortex at...
each contact location, the maximum d.c. current \( I_{\text{dc}} = 40 \text{ mA} \) was applied. The spectra were then recorded at decreasing d.c. current values. The plots at fixed \( I_{\text{dc}} \) (see Fig. 2b for \( I_{\text{dc}} = 24 \text{ mA} \)) allow interpretation of the device output as the superposition of four signals generated by four independent (or weakly interacting) vortices orbiting around the contacts. Four peaks are seen to maintain their relative frequency positions through the whole current range. The small peak frequency differences are due to the unavoidable and uncorrelated differences in the nanocontacts size, which lead to slightly different currents flowing through the four contacts connected in parallel. Similar results have been obtained (not shown) for the 10-μm-spaced contacts.

The frequency increases with d.c. current (blueshift) with a tunability of \( df/dI_{\text{dc}} = +15.7 \text{ MHz mA}^{-1} \) and \( df/dI_{\text{dc}} = +19.6 \text{ MHz mA}^{-1} \) for the devices with \( d = 2 \mu\text{m} \) and \( d = 10 \mu\text{m} \), respectively. At a fixed d.c. current, we observed an increase of the frequency with perpendicular applied field at a rate of \( df/d\mu_0 H_z = +2.5 \text{ MHz mT}^{-1} \) and \( df/d\mu_0 H_z = +2.3 \text{ MHz mT}^{-1} \) for the samples with \( d = 2 \mu\text{m} \) and \( d = 10 \mu\text{m} \), respectively. Note that, for both the devices with \( d = 2 \mu\text{m} \) and \( d = 10 \mu\text{m} \), no signal was detected at \( H_z = 0 \text{ mT} \), but we had to apply a few tens of mT with the same polarity as the reset field to sustain radio-frequency output. The total integrated power, numerically estimated as the area under the acquired spectrum, is \( 22 \pm 2 \text{ nW A}^{-2} \) in the case of intercontact spacing \( d = 2 \mu\text{m} \) and \( 23 \pm 2 \text{ nW A}^{-2} \) for \( d = 10 \mu\text{m} \). The power is independent of \( I_{\text{dc}} \) and \( H_z \) in the studied range and not significantly dependent on \( d \).

**Phase-locking of closely spaced nano-oscillators**

The spectrum map of the device with \( d = 500 \text{ nm} \) (Fig. 2c) is much richer. It shows a knee below which four peaks can again be resolved. At larger bias currents, the four peaks merge into two peaks, over a small range of currents, and then into a single peak (Fig. 2d). The merging is accompanied by a sudden increase of the power amplitude and a significant narrowing of the linewidth (narrowest measured linewidth \( \Delta f = 0.9 \text{ MHz} \)). Let us emphasize that the spectrum map shown in Fig. 2c was recorded in zero applied field. We define the critical current \( I_{\text{cr}} \) as the smallest current at which a single peak can be resolved in the spectrum. In zero field, \( I_{\text{cr}} = 22.0 \pm 0.5 \text{ mA} \) and it changes with the applied field at a rate \( dI_{\text{cr}}/d\mu_0 H_z = +0.1 \text{ mA mT}^{-1} \).

The output was sustained in a field range around zero, typically for \( |H_z| \) smaller than 10 mT (Fig. 2e and f), that is, in a range where gyrotropic precession was not observed in widely spaced (\( d = 2 \) and 10 µm) nanocontacts. The largest peak amplitude was recorded at \( H = 0.0 \text{ mT} \). An applied field always leads to a reduction of the peak amplitude and an increase of the linewidth, with the area underneath remaining constant, regardless of the sign of the field. The range of oscillating frequency is significantly lower than the one found for the widely spaced nanocontacts (see the range of oscillation frequencies in Fig. 2a and c). In fact, a frequency (in the centre of the four peaks) \( f = 581.5 \text{ MHz} \) at \( I_{\text{dc}} = 25 \text{ mA} \) and \( H_z = 50 \text{ mT} \) for the case \( d = 2 \mu\text{m} \) should correspond, according to the measured tunability in field, to a frequency \( f = 456.5 \text{ MHz} \) at \( H_z = 0 \text{ mT} \) and at the same \( I_{\text{dc}} \). This value is about twice the oscillating frequency in the closely spaced (\( d = 500 \text{ nm} \)) nanocontacts at the same \( I_{\text{dc}} \). Following this frequency behaviour, we also find that the tunability in current is reduced in the closely spaced device compared to the others to \( df/dI_{\text{dc}} = +5.6 \text{ MHz mA}^{-1} \).

From the spectra in Fig. 2d, we estimated an integrated power \( P = 88 \pm 2 \text{ nW A}^{-2} \) for \( I_{\text{dc}} > I_{\text{cr}} \). Just below \( I_{\text{cr}} \), the total integrated power decreases suddenly to reach half the power at \( I_{\text{dc}} = 18.5 \text{ mA} \) and keeps decreasing. This behaviour is similar to that observed for the unlocking of uniform modes in arrays of nanocontacts coupled through spin waves\(^{12,13}\). At much lower \( I_{\text{dc}} \), where the remaining
The starting point of our analysis is the determination of the ground state. At room temperature, the maximum intercontact distance at which this condition is satisfied can be estimated by comparing the far-field dipolar energy of the system with the thermal energy. The dipolar energy in a vortex–antivortex pair can be estimated in the following way. In a vortex (antivortex) the out-of-plane magnetization is significantly different from zero only in the core region. Thus, over distances much larger than the core size, the stray field can be written as the field generated by a dipole of magnetic moment $\mu$ (refs 17, 19).

$$U_{\text{dip}} \approx \frac{\mu^2}{\delta l}$$

with $\delta$ being the vortex–antivortex distance. The associated characteristic temperature $T_{\text{dip}} = U_{\text{dip}}/k_B$, with $k_B$ being the Boltzmann constant, is equal to room temperature for $\delta \approx 215$ nm. This estimation is consistent with our observation of phase-locking at room temperature for intercontact distance $d = 500$ nm.

Although the antiferromagnetic chessboard is stabilized by the dipolar interaction between vortices and antivortices, this interaction, for the rather large intercontact distances considered here, is not strong enough to keep the vortices phase-locked. In fact, the long-range exchange interaction between a vortex and the nearest-neighbour antivortex is 16.

$$U_{\text{ex}} = AL \ln \frac{\delta}{\nu}$$

where $A$ is the exchange constant, $\nu = 2l_{\text{ex}}$ is the size of the vortex core, and $\delta$ is assumed to be equal to 250 nm. The exchange energy is $U_{\text{ex}} = 1.4 \times 10^{-18}$ J, which is about three orders of magnitude larger than the long-range dipolar energy estimated through equation (2) at the same distance. This means that the vortex coupling is due to the balance between the restoring force exerted by the periodic Oersted potential and the attractive forces between vortices and antivortices due to the exchange interaction.

The exchange interaction between the vortices and their surrounding antivortices forces the former to precess on a larger orbit, with a consequent reduction of the detected output frequency and tunability in current. This reduction does not seem to be directly related to the phase-coupling but simply to the reduction of $d$, and, hence, of the vortex–antivortex distance. In fact, a reduction of frequency and tunability is also observed when decreasing the intercontact distance from $d = 10 \mu$m to $2 \mu$m.

The vortex–antivortex stray field at a distance $\delta = 250$ nm is rather small, that is, $\sim 0.01$ mT calculated as $U_{\text{dip}}/\mu$. Therefore, in principle, any perpendicular field larger than this, applied during
nucleation, should inhibit the formation of the antiferromagnetic chessboard. In contrast, we have already shown in Fig. 2f that condensation occurs in a much larger range of $H_{\perp}$.

The behaviour in the presence of a perpendicular applied field can be understood with the help of a dynamic micromagnetic simulation. When a field is applied perpendicular to the plane during nucleation of the vortices, the static ground state is one of the ferrimagnetic configurations\(^\text{37}\). For instance, in a simplified system of two contacts, if a field of $+1$ mT is applied, all the vortices and antivortices nucleate with the core polarities oriented as the field (see Supplementary Video 1). During their incoherent oscillation, the distance between vortices and anti-vortices periodically falls below $d/2$, causing an increase in the energy of the system. In a certain range of externally perpendicular applied d.c. fields, this distance can be sufficiently small that the energy of the ferrimagnetic configuration is dynamically larger than that of the antiferromagnetic chessboard. As a consequence, the system goes into the antiferromagnetic state by switching of the antivortex core. In the experiment, the transient occurs on a timescale much smaller than the acquisition time, and hence is not recordable during spectrum acquisition.

This switching mechanism has some similarities with the recently discovered switching of a vortex core through bursts of a.c. field\(^\text{20,21}\) or a.c. current\(^\text{21}\). In references 20 and 21, the authors demonstrate that the topology of an orbiting vortex changes when its velocity approaches a critical value. Regions with magnetization opposite to that of the core polarity appear. In these conditions, the field amplitude needed to switch the vortex core is strongly reduced compared to the case of a static vortex. In our system, the stray field of the orbiting vortex indeed acts as an a.c. local field on the antivortex core (and vice versa). However, our simulations suggest that the topological change of the antivortex structure, which allows its core switching with the small field bursts provided by the orbiting vortices, is caused by the exchange interaction with the approaching vortices rather than by dynamic self-deformation when reaching a critical velocity.

In the presence of an in-plane field during nucleation, the spectrum map becomes very rich, as shown in Fig. 4a. The observed multiple peaks are clearly correlated. For instance, at $I_{\text{dc}} = 27.5$ mA, where four peaks can be observed (Fig. 4b), they are symmetric around the oscillation frequency of the locked mode in the same field conditions with the system previously stabilized in zero in-plane field. The total integrated power is $83 \pm 2$ nW A\(^{-2}\), that is equal to that measured in the locked state and not a quarter of it, as would be the case if the spectrum were simply the superposition of the signals generated by four vortices independently orbiting around the contacts. Moreover, we emphasize that for other current values (see Fig. 4a) or in-plane applied fields (see Supplementary figure), a different number of modes appear in the spectrum. This implies that the number of excited eigenmodes is not related to the number of nanocontacts.

The effect of an in-plane field can be understood once again with the help of a dynamic micromagnetic simulation. The in-plane field $H_{\|}$ elongates the orbits and off-centres them along the plane direction perpendicular to $H_{\|}$ (see Supplementary Video 2). This is equivalent to symmetry breaking. By making the two-dimensional potential asymmetric, the strong nonlinear system can be driven from a to symmetry breaking. By making the two-dimensional potential field conditions with the system previously stabilized in zero in-plane field (see Supplementary Video 1). During their incoherent oscillation, the distance between vortices and anti-vortices periodically falls below $d/2$, causing an increase in the energy of the system. In a certain range of externally perpendicular applied d.c. fields, this distance can be sufficiently small that the energy of the ferrimagnetic configuration is dynamically larger than that of the antiferromagnetic chessboard. As a consequence, the system goes into the antiferromagnetic state by switching of the antivortex core. In the experiment, the transient occurs on a timescale much smaller than the acquisition time, and hence is not recordable during spectrum acquisition.

Dynamics of interacting solitons, a laboratory system

Let us finally point out that, if one considers a matrix with a much larger number of contacts, it has been theoretically proved\(^\text{16}\) that our system is formally equivalent to a variety of two-dimensional systems, such as superfluid films and two-dimensional magnets. In fact, a transition to the superfluid phase occurs when pairs of vortex and antivortex appear, and eventually condensate into a vortex–antivortex square lattice. The transition is called the Berezinskii–Kosterlitz–Thouless (BKT) transition\(^\text{24}\) and can be regarded as the two-dimensional counterpart of a Bose–Einstein condensation (BEC)\(^\text{25,26}\). BKT (as well as BEC) transitions are experimentally observed when two conditions are satisfied: the solitons are prevented from spreading by attracting them into potential wells, and at the same time, the solitons are prevented from shrinking by stirring the condensate. The vortices generated by the d.c. current in our array of nanocontacts are prevented from spreading by the force due to the periodic Oersted potential. Because the d.c. current is spin-polarized by the second ferrimagnetic layer, the spin transfer torque provides the stirring force on the vortices. The theory of two-dimensional solitons strictly applies because the vortices precess in a continuous film\(^\text{27}\). If the precessing vortices interact through antivortices, phase-locking can be used to model a BKT transition.

Methods

For the devices preparation, the multilayer tantalum (5 nm)/copper (40 nm)/permalloy (4 nm)/gold (6 nm) was deposited on a silicon substrate by d.c. magnetron sputtering (without breaking vacuum) in an argon atmosphere. A resist was spun on the sample for subsequent conductive tip AFM indentation. The resist thickness was tuned to be 40 nm by diluting it with solvent. A voltage of 4 V was applied between the tip and the sample during indentation. By monitoring the tip–sample resistance, the indentation could be stopped when the distance between the tip and the surface of the sample was a few nanometres. The sample was then transferred back in the vacuum chamber where the holes were enlarged and the sample surface reached by etching the resist in 15 W O\(_2\) plasma for 45 s.
breaking vacuum, the cobalt (15 nm) / gold (100 nm) top electrode was deposited. Finally, the top electrode was defined by ion milling with real-time monitoring by mass spectroscopy.

For high-frequency measurements, the total direct current $I_{ex}$ was applied to the contact matrix through a bias tee. The excited oscillation voltage was amplified (25 dB) and recorded with a commercial spectrum analyser. The shown power spectra were extracted from those recorded by, subsequently, subtracting a reference spectrum measured at $I_{ex} = 0$ mA in the same magnetic field conditions and removing the amplification.

A field of 0.7 T was applied out of the sample plane as a reset field before acquisition of each spectrum. The polarity of this field defines the positive field direction. In all the shown spectra, the magnetic field is always applied before acquisition of each spectrum. The polarity of this field defines the positive field direction. In all the shown spectra, the magnetic field is always applied before applying the current.

The micromagnetic simulations were performed with the public code OOMMF. The simulations included the Oersted field due to the current. For calculation of this field, we assumed the infinite wire approximation with current density constant in the field, we assumed the infinite wire approximation with current density constant in the area of the contacts. A cell size of 2.5 nm $\times$ 2.5 nm $\times$ 5 nm was used for the simulation shown in Fig. 3 and Supplementary Video 1, and 5 nm $\times$ 5 nm $\times$ 5 nm for Supplementary Video I. We used the following material parameters for cobalt: saturation magnetization $M_S = 1.4 \times 10^6$ A m$^{-1}$, exchange constant $A = 3 \times 10^{-11}$ J m$^{-1}$. Thus, the exchange length was calculated to be $l_{ex} = \sqrt{2A/\mu_0M_S} = 5$ nm. A damping constant $\alpha = 0.01$ and a polarization $P = 0.3$ were used in the dynamic simulations.

Received 27 March 2009; accepted 5 May 2009; published online 21 June 2009

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