Self organized mode locking effect in superconductor / ferromagnet hybrids

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The vortex dynamics in a low temperature superconductor deposited on top of a rectangular array of micrometer size permalloy triangles is investigated experimentally. The rectangular unit cell is such that neighboring triangles physically touch each other along one direction. This design stabilizes remanent states which differ from the magnetic vortex state typical of individual non-interacting triangles. Magnetic Force Microscopy images have revealed that the magnetic landscape of the template can be switched to an ordered configuration after magnetizing the sample with an in-plane field. The ordered phase exhibits a broad flux flow regime with relatively low critical current and a highly anisotropic response. This behavior is caused by the spontaneous formation of two separated rows of vortices and antivortices along each line of connected triangles. The existence of a clear flux flow regime even for zero external field supports this interpretation. The density of induced vortex-antivortex pairs is directly obtained using a high frequency measurement technique which allows us to resolve the discrete motion of vortices. Strikingly, the presence of vortex-antivortex rows gives rise to a self organized synchronized motion of vortices which manifests itself as field independent Shapiro steps in the current-voltage characteristics.

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I. INTRODUCTION

The dynamics and pinning of vortices in type-II superconductors is arguably one of the most widely investigated phenomena in the field of superconductivity. Recently, particular attention has been devoted to the possibility of controlling the strength of the vortex pinning using different magnetic templates. Within the London formalism, the interaction between a vortex line and a finite size permanent magnet is given by

\[ U_p(R) = -\int_{\text{dot}} m(r) \cdot B_v(R-r) \, d^3r, \quad (1) \]

where the integration is carried out over the volume of the ferromagnet, \( R \) indicates the position of the vortex line, \( m(r) \) its spatial dependent magnetic moment, and \( B_v \) is the field generated by the vortex line. This equation shows that the pinning potential not only depends on the size of the dot but also on their exact magnetic state. In particular, if the superconducting penetration depth \( \lambda \) is on the order of the lateral size of the dot, a weak average pinning is expected when the magnetic dot contains multiple domains. In contrast to that, a maximum pinning should be achieved for single domain structures. In order to retain these properties in big arrays of magnetic particles it is necessary to ensure a minimum distance between them in such a way that the field profiles of neighboring elements do not overlap and can be resolved at scales of the magnetic penetration length \( \lambda \).

Notice that according to Eq. (1) it is possible to increase strongly the pinning by just growing continuously the magnetic moment of the dots. This growth is limited by the fact that above certain magnetic moment, a vortex-antivortex pair will be induced by the magnetic element with the consequent change in the pinning properties. For instance, dots with perpendicular magnetic moment induce a vortex on top of the dots whereas the accompanying antivortex is located at interstitial positions, in between the dots. As a result, vortices and antivortices feel a different pinning potential which leads to a distinct dynamic response. Since the weakly pinned interstitial antivortices limit the maximum critical current \( J_c^{\text{max}} \), their annihilation with an external field produces a shift of the field position of the \( J_c^{\text{max}} \) towards \( H = nH_1 \), where \( n \) is the number of vortex-antivortex pairs generated by a single magnetic element and \( H_1 \) is the field corresponding to one external vortex per unit cell. Since \( H_1 \) is a geometrically predefined parameter, the field loci of the \( J_c^{\text{max}} \) is a direct measurement of the number \( n \) of bounded pairs. Unfortunately, in the case of micromagnets with in-plane magnetic moment, the field polarity symmetry of the system conceals this information from dc electro-transport experiments.

In this work, we introduce alternative methods to accurately determine the number \( n \) of vortex-antivortex pairs induced by the magnetic template, based on the ac-dynamic response of the hybrid system. Indeed, the
fact that the critical force to bring the vortex and antivortex closer differs from that necessary to pull them apart, gives rise to a clear rectification effect when the system is submitted to a zero average ac-signal. This rectified voltage is maximum at zero external field and progressively decreases with increasing field until all vortices (or antivortices) are compensated. At RF frequencies the mode locking or synchronized motion of vortices produces distinct voltage steps in the current-voltage characteristics with separation determined by the number of moving vortices. Both techniques indicate that for the particular system here investigated, the magnetic template generates one single vortex-antivortex pair per unit cell.

The crucial conditions for the success of these methods are (i) a periodic landscape, i.e., all vortices and antivortices feeling the same environment, and (ii) a free flux flow motion of vortices. Although the predefined topological landscape is by definition periodic, the magnetic landscape for the as-grown state consists of a highly disorder distribution of magnetic poles. This disordered state can be switched to a perfectly ordered magnetic lattice by magnetizing the sample with an in-plane field. The order-disorder transition, already reported for other S/F systems, produces a profound modification in the vortex dynamics. More specifically, in the disordered state, non-linear current-voltage characteristics with high critical currents are observed, whereas a broad flux flow regime accompanied by a sharp decrease of the critical currents is obtained in the ordered state. This behavior can be attributed to the reduction of the average hopping distance determined by the separation between neighboring magnetic poles of identical polarity. In addition, the presence of a flux flow regime at zero applied field reinforces the idea that the investigated magnetic structure is able to induce vortex-antivortex (v-av) pairs.

II. SAMPLE PREPARATION AND CHARACTERIZATION

The sample investigated consists of a periodic array of permalloy (Ni$_{80}$Fe$_{20}$) equilateral triangles of 1 µm side and a thickness of 25 nm. Along one of the principal directions of the array (x-axis) the separation between the triangles is 1 µm whereas in the perpendicular direction (y-axis) the triangles are touching [see Fig. 1(a)]. This geometry gives rise to two well defined directions of the structure, parallel (y) and perpendicular (x) to the line connecting nearby triangles, as indicated in Fig. 1(a). A 50 nm thick superconducting Al bridge deposited on top of the magnetic template also covers the unpatterned substrate [see Fig. 1(b)] thus allowing us to compare directly patterned vs reference plain film and identify the influence of the magnetic triangles on the superconductor. Both, the array of triangles and the transport bridge are aligned at submicrometer scale and were prepared by standard lithographic techniques. The Al plain film has a critical temperature at zero field $T_{c0} = 1.315$ K, with residual resistivity ratio $R(300K)/R(4.2K) \approx 2$, and superconducting coherence length $\xi(0) \approx 128$ nm.

It is a crucial point that neighboring triangles overlap in one direction in order to eventually achieve two well distinguished phases: disordered in the as-grown state and ordered in the magnetized state. Indeed, Magnetic Force Microscopy (MFM) images shown in Fig. 1(c) for non touching triangles demonstrate that the remanent state, irrespective of the magnetic history of the triangles, is always a vortex state, in agreement with previous reports on individual micrometer size magnetic structures. In contrast to that, if there is a physical contact between neighboring triangles a new set of possible states can be found. Fig. 1(d) shows a MFM image of the sample investigated in the as-grown state. Here a clearly disordered arrangement of the, so called, buckle states is observed. Interestingly, after magnetizing the sample with an homogeneous in-plane field perpendicular to the chain of triangles [Fig. 1(e)], these buckle states can be set in a perfectly ordered arrangement. The in-plane distribution of magnetic moments calculated by OOMMF micromagnetic simulations is indicated in Fig. 1(f) with black arrows whereas the z-component of the stray field ($B_z$) is shown in Fig. 1(g). In all cases the magnetic state of each triangle is characterized by a strong out of plane magnetic pole which is located at one of the two bottom corners of the Py triangles and a weaker magnetic pole, with opposite polarity, located at the top corner of the triangle [Fig. 1(e)].

![Image](image.png)

FIG. 1: (color online) (a) schematic drawing of the array of Py triangles. The direction of the external field, the current, and the relative orientation of the array is clearly indicated with arrows. (b) Schematic representation of the Al transport bridge with Hall voltage contacts on the patterned region. Magnetic Force Microscopy pictures obtained at remanence of physically separated magnetic triangles (c), touching magnetic triangles in the virgin state (d), and the triangles magnetized perpendicular to the chain of triangles (e). Panels (f) and (g) show the local magnetic moment distribution and the z-component of the stray field, respectively, calculated by OOMMF simulations.
is precisely this magnetic pole which can be controlled to lie on either side of the triangle by simply changing the direction of magnetization. Notice that this magnetic monopole-like structure breaks the field polarity symmetry of the S/F hybrid system. The question now arises as whether the distribution of the magnetic dipoles (ordered versus disordered) has any effect on the properties of the adjacent superconducting film.

III. NUCLEATION OF SUPERCONDUCTIVITY

Let us first analyze the nucleation of the superconducting order parameter for the two different magnetic states of the hybrid system. Fig. 2 summarizes the obtained superconductor/normal metal phase boundaries for a dc-current of 100 µA and using 90% criterion of the normal state resistance $R_n$ for different magnetic states and current orientations. Higher criteria exhibit similar results. For the as-grown state (open circles) a pronounced decrease of the critical temperature at low fields with respect to the plain film is observed. A further suppression of $T_c$ is obtained after the sample has been magnetized with an in-plane field parallel to the $x$ direction (solid triangles). The fact that the phase boundary of the as-grown state lies in between those corresponding to the plain film and the fully polarized buckle state, suggests that in the virgin sample, randomly oriented buckle states coexist with magnetic vortex states. This has been indeed confirmed by the MFM images taken at different locations of the pattern.

It is worth noticing that for the magnetized sample the $T_c(H)$ boundary exhibits a parabolic background similar to that reported in earlier investigations. This behavior can be ascribed to the change of dimensionality when $\xi(T)$ exceeds the size $w$ of the regions where superconductivity first nucleates. Within this regime the phase boundary can be approximated by $T_c(H)/T_c(0) = 1 - (\alpha H)^2$, with $\alpha = \xi(0)\pi w/\sqrt{3}\phi_0$. By using this expression to fit the data in Fig. 2 (solid line) we found $w \sim 1.4 \, \mu m$ which is consistent with the largest separation between neighboring magnetic poles. This behavior suggests that for the magnetized sample the stray fields confine the superconducting condensate in between the magnetic poles and consequently the superconducting order parameter will be modulated into a sequence of quasi one-dimensional stripes. This prediction is in agreement with complementary transport measurements for two orthogonal orientations of the applied current, perpendicular to the chain of triangles (solid triangular symbols) and parallel to them (open triangular symbols) as shown in Fig. 2 indicating that the onset of superconductivity is independent on the direction of the applied current.

IV. VORTEX DYNAMICS IN THE ORDERED AND DISORDERED PHASES

In the previous section we showed that by switching from the as-grown (disordered phase) to the magnetized (ordered phase) state, a proliferation of magnetic poles takes place which in turn gives rise to a decrease of $T_c$. In addition, the dimensional crossover from a 2D disordered phase to a 1D ordered phase has little or no influence on the nucleation temperature. However, deeper into the superconducting state where the electric response is dominated by the vortex dynamics, the one dimensional nature of the field modulation should become more apparent.

In order to address this issue we have measured the voltage-current characteristics $[V(I)]$ with the bias current applied in the $x$-direction at $H = \frac{1}{2}H_1$ and $T = 1.1$ K for the two different magnetic states: as-grown (open circles) and magnetized (solid squares) states as shown in Fig. 3 (a). The most obvious feature in this Figure is the enormous difference (more than an order of magnitude) between the critical current of the magnetized state (24 µA) with respect to the demagnetized state (490 µA). Furthermore, the vortex dynamics is severely modified after magnetizing the sample. Indeed, whereas in the as-grown state there is a narrow window of currents exhibiting a highly non-linear $V(I)$ dependence, where the vortex lattice transit from pinned phase to the normal state, in the magnetized sample a broad flux-flow...
regime is observed immediately above the critical current. Fig. 3(b) and (c) summarize the field dependence of the different dynamical phases for the as-grown sample [panel (b)] and after magnetization [panel (c)]. The onset of vortex motion is delimited by a resistance criterion of 10% \( R_n \) while a 90% \( R_n \) criterion is used to indicate the transition to the normal state.

![Fig. 3](image.png)

FIG. 3: (color online) (a) Voltage-current \( V(I) \) characteristic with current applied in the \( x \)-direction, at \( H = \frac{1}{2} H_1 \) and a temperature of 1.1 K for the as-grown and the magnetized states. The lower panels show the magnetic field dependence of the critical current and the transition to the normal state as determined by criterion of 10% \( R_n \) and 90% \( R_n \), respectively, for the as-grown (b) and magnetized (c) states.

Notice that the presence of matching features, which require two-dimensional phase coherence, together with the fact that in the magnetized state a finite critical current is needed to depin the vortices, rule out the possibility of having normal metal stripes in the hybrid sample. Instead, the observed difference between both phases can be attributed to the formation of easily moving channels of vortices and antivortices created by the magnetic structure. Indeed, in the ordered phase the potential wells created by the magnetic dipoles are perfectly aligned and strongly overlap, thus creating an easy channel for flux motion. In contrast to that, in the disordered phase, the average distance between two neighboring magnetic pinning centers of equal polarity is larger thus impeding the easy motion of vortices. From now on we will focus on the dynamic response of the system in the ordered state.

V. EVIDENCE OF INDUCED V-AV PAIRS

A rough criterion to determine whether the stray field emanating from the magnetic structure is strong enough to induce v-av pairs can be obtained by comparing it to the penetration field \( H_p \) resulting from Bean-Livingston barriers.\(^{16}\) The highest barrier possible is limited by the thermodynamic critical field \( H_p = \phi_0/(2\sqrt{2}\pi\lambda\xi) \). Using \( \xi = 128 \) nm and \( \lambda = 150 \) nm we obtain \( H_p \sim 12 \) mT which lies well below the field intensity (\( \sim 100 \) mT) produced by the magnetic template at the superconducting film. Based on this estimation we can safely conclude that the magnetic structure is able to induce, at least, one vortex-antivortex per unit cell.

![Fig. 4](image2.png)

FIG. 4: (color online) Schematic drawing of the induced vortex lattice in the case of no external applied field (a) and an applied field equal to \( \frac{1}{2} H_1 \) (b). Red (green) circles represent positive (negative) vortices.

This finding suggests that the ordered phase consists of two well defined channels with opposite vortex polarity as schematically shown in Fig. 4(a). Since the magnetic template breaks the field polarity symmetry, in principle vortices and antivortices should experience a different depinning current. Moreover, as shown in Fig. 4(b), under these circumstances an external field will compensate vortices in the channel with opposite polarity (thus introducing disorder in this particular channel and increasing the average hopping distance) and as a result the high vortex mobility will be suppressed. In contrast to that, the row with vortices of opposite polarity would remain unaltered. This scenario predicts two anomalous behaviors: the existence of a flux-flow response at zero field and a field independent flux flow regime as long as none of the two rows of vortices is fully compensated by the external field. Both predictions are in agreement with the field evolution of the \( V(I) \) measurements done.
at 1.1 K. The slope of the linear regime, proportional to the amount of particles contributing to the easy flow, is plotted as a function of field in Fig. 5(a). In comparison to the slope expected for the free flux flow model $R_{ff} = R_n H/H_{c2}$ [straight line in Fig. 5(a)], where $R_n$ is the normal state resistance and $H_{c2}$ is the upper critical field, a weak field dependence of the flux flow slope is seen at low field values.

Further evidence in favor of the existence of channels populated with v-av pairs comes from the measurements of the critical current for currents applied parallel ($J_y$) and perpendicular ($J_x$) to the channels of v-av pairs. These results are shown in Fig. 5(b), with a critical current determined using a criterion of 100 $\mu$V. This figure shows a strong in-plane anisotropy $\gamma = J_x/J_y$ induced by the magnetic template which results from the easier motion of vortices along the line of magnetic poles than perpendicular to them. This anisotropy amounts to a factor of 5 at non matching fields and is minimum at the commensurability fields.

VI. DETERMINATION OF THE V-A V PAIRS DENSITY AT LOW FREQUENCIES: RATCHETS

The scenario described in the previous section where rows of vortices and antivortices generated by the magnetic template coexist, makes our system similar to the array of elongated magnetic bars studied by de Souza Silva et al. This system has been modeled within the London approximation by Carneiro for the case where no v-av pairs are induced, and more recently by Lima and de Souza Silva assuming that v-av pairs emanate from the dipoles. Irrespective of whether the magnetic template is able to induce v-av pairs or not, both situations lead to a rectification of the net vortex motion $v$ when a zero average ac-force is applied parallel to the dipolar moment (i.e. along the line connecting a vortex with its corresponding antivortex). However, clear differences in the field dependence of the average velocity are expected. Indeed, if v-av pairs are present then (i) a maximum rectification is expected at zero field, (ii) unlike standard ratchet signal, the sign of the average voltage readout should be field polarity independent, (iii) as vortices (antivortices) are progressively compensated by increasing the external negative (positive) field, the rectified signal should decrease.

The measured rectification effect, shown in Fig. 6, confirms all the above mentioned facts and indicates that the v-av pairs are more easily brought together than separated. Furthermore, if the applied field is increased above the first matching field a sign reversal occurs. This sign reversal at $H_1$ can not be explained assuming more than one v-av pair created by the magnetic template since in this case the v-av annihilation process would be unlikely interrupted by the partial compensation of one row of vortices. In contrast to that, if there is a single v-av pair per unit cell, once a whole line is compensated by the external field, a dramatic change in the rectification properties is expected. Although this finding seems to indicate that the magnetic structure is able to induce a single v-av per triangular element, a more convincing evidence of this prediction can be obtained only by monitoring the individual motion of vortices at higher frequencies.
VII. DETERMINATION OF THE V-AV PAIR DENSITY AT HIGH FREQUENCIES: MODE LOCKING

The motion of vortices in a periodic potential gives rise to a variety of interesting dynamical effects such as the generation of rf radiation and synchronization/resonance effects of the vortex lattice when a rf current \( I_{rf} \) is superimposed onto a dc drive \( I_{dc} \). The latter effect gives rise to a series of plateaus in the \( V(I) \) measurements very similar to the Shapiro steps observed in Josephson junctions. The relevant timescale of these processes is determined by the washboard frequency, \( f = v/d \) with \( d \) the period of the potential and \( v \) the average velocity of the particles. The investigation of the above mentioned phenomena can be used to obtain additional information on the dynamics of the vortex lattice, the periodicity of the underlying potential and the density of particles in the system. In other words, high frequency experiments provide further insight into the microscopic motion of vortices not available with the averaged response presented in previous sections.

The synchronous motion of vortices resulting from the resonant condition \( f = v/d \), known as mode locking effect, allows one to obtain the exact number of v-av pairs created by the magnetic template. Indeed, the value of the average velocity is given by:

\[
v = cf d
\]

with \( d \) the periodicity of the lattice in the direction of motion, \( f \) the frequency of the rf current and \( c \) the number of cells that a vortex has traveled during a rf cycle. As we mentioned above, the fingerprint of this type of resonant motion is the presence of clear plateaus in the \( I(V) \) characteristics separated by voltage steps,

\[
V_c = Nk \frac{h}{2e} \frac{v}{d} = Nk \frac{h}{2e} cf,
\]

with \( k \) the number of vortices in each unit cell and \( N \) the number of rows between two voltage contacts. Measuring the voltage separation between consecutive steps and using Eq. (3) we can estimate the density of vortices moving coherently in the flux flow regime.

Fig. 7(a) shows a \( V(I) \) scan for \( H = 0.5 \text{ mT} \) and \( T = 1.1 \text{ K} \) with a small ac excitation \( I_{rf} \ll I_{dc} \) of 100 MHz superimposed. The presence of Shapiro steps in the linear part of the \( V(I) \) curve becomes more apparent by plotting \( V \) versus \( dI/dV \) [see Fig. 7(b)]. These steps represent a clear indication of the coherent flow of vortices.

In Fig. 7(b) we assign the present peaks to values of \( V_c \) with \( c \) ranging from 1 to 4 (\( V_1 = 63 \mu V \), \( V_2 = 127 \mu V \), \( V_3 = 194 \mu V \) and \( V_4 = 261 \mu V \)). From their linear dependence on \( c \), shown in the inset of Fig. 8 we calculate a value of \( k \approx 1.06 \). An independent estimation can be obtained by tracking the frequency dependence of each voltage plateau \( V_c \). This is shown in Fig. 8 for the first two steps. The linear fit using Eq. (3) gives \( k \approx 1.1 \) and 1.18 for \( V_1 \) and \( V_2 \), respectively.

It is important to notice that, in general, two ingredients are necessary for the presence of mode locking effects. Firstly, the existence of a flux flow regime, which ensures a narrow distribution of average vortex velocities. Secondly, there should be a periodic perturbation of the pinning potential, otherwise the distortion of the vortex lattice becomes too strong and its coherent motion is destroyed. In principle, this second condition is only fulfilled at the matching field, \( H_1 = 1.195 \text{ mT} \). This seems to be in contradiction with the fact that clear Shapiro steps are present at 0.5 mT, a field which would correspond to 0.42 vortices per unit cell. This apparent paradox can be solved invoking once again the presence of v-av pairs induced by the triangles. In this case, the applied field will annihilate the vortices in the channel with opposite polarity, leading to an increase of their critical current and consequently suppressing their contribution to the coherent motion. On the other hand, the row with equal polarity will be unaltered by the applied magnetic field and therefore should exhibit a region of field independent Shapiro steps. In other words, the incoming vortices generated by the external magnetic field play no role in the coherent oscillatory motion of vortices thus giving rise to a constant voltage separation between two consecutive Shapiro steps.

In order to check this hypotheses of self organized commensurability we measured the field dependence of the Shapiro steps, which are better visualized as a peak in the \( dI/dV \) as previously shown in Fig. 7(b). The results of these measurements for \( T = 1.1 \text{ K} \) are summarized in Fig. 9 as a contour plot in the field-voltage plane, with
maximum intensity representing the exact location of the $V_c$ peaks. For magnetic field values ranging from $-H_1$ to $+H_1$ a clear field independence peak position is observed. This result confirms that the amount of particles, moving coherently, is independent of the applied external field.

**FIG. 8:** Frequency dependence of $V_1$ (circular symbols) and $V_2$ (square symbols) at a field of 0.5 mT and temperature of 1.1 K. Inset: $V_1 \rightarrow V_4$ measured at 100 MHz.

**FIG. 9:** (color online) Magnetic field dependence of $dI/dV$ versus $V$ at a temperature of 1.1 K and an applied rf frequency of 100 MHz.

**VIII. DISCRETE MOTION OF THE VORTEX-ANTIVORTEX LATTICE**

It has been predicted by molecular dynamics simulations that a vortex lattice in a periodic landscape of asymmetric pinning sites can unveil the discrete hopping of vortices when the system is excited with a zero mean rf drive. So far, most of the experimental reports evidence a vortex rectification effect obtained in the low frequencies limit which results from an averaging in time and are not able to resolve the individual motion. In this section, we will try to bridge this gap by constructing an artificial vortex rectifier resulting from a dc tilt, which breaks the inversion symmetry of our system. Additionally, this tilt allows us to probe the discrete motion of the vortex lattice without the need of a time resolved signal read out.

**FIG. 10:** DC voltage as a function of ac amplitude ($f = 100$ MHz) with a dc offset of 10 $\mu$A applied. The indexing $(a, b)$ indicates the number of steps forward $(a)$ and backward $(b)$.

Fig. 10 shows the dc voltage as a function of the applied ac drive with an additional dc current on top ($I_{dc} = 10$ $\mu$A). Due to the applied dc tilt the critical current, needed to depin vortices, is smaller for positive than for negative currents ($I_{pos} < I_{neg}$). Depending on the strength of the applied ac amplitude $I_{ac}$ we can determine three different dynamic regimes. For small ac excitations ($I_{ac} < I_{pos} < I_{neg}$) the driving force is unable to depin the vortices and as a result the dc voltage is zero. By increasing the ac amplitude above $I_{pos}$ the vortex lattice is set in the flux flow regime and mode locking effects result in an oscillatory dependence of the dc voltage as a function of the ac amplitude. Finally at high ac amplitudes the normal state is reached and $V_{dc} = R_n I_{dc}$. Each maxima and minima in dc voltage can be attributed to the discrete motion of the vortex lattice in the positive and negative cycle of the driving amplitude. To identify the motion, we introduce the indexing $(a, b)$ in Fig. 11 with $a$ and $b$ the number of steps in the positive and negative direction respectively. At $I_{ac}$ just above $I_{pos}$ the driving force is strong enough to move the vortex lattice one unit cell in the positive direction while no motion occurs in the negative cycle of the ac drive and a maximum in the dc voltage output occurs. If the ac amplitude passes $I_{neg}$ a single back and forth step of the vortex lattice occurs and as such a strong reduction of the dc voltage is measured. This difference between back and forth motion of the vortex lattice, due to a small dc tilt, results in a discrete stepping procedure up to the sixth order.

Fig. 11 shows the dependence of the dc voltage, normalized to $V_1$, as a function of the ac drive and the dc tilt. At small dc tilts (below $\approx 100 \mu$A) the behavior is exactly the same as shown in Fig. 10. In this region, the applied dc current creates a small asymmetry and
the vortex lattice, depending on the exact value of the ac drive, can either move one step extra in the positive direction (maxima in dc voltage) or an equal amount of steps in both directions (minima in dc voltage). Around 100 $\mu$A the oscillations vanish (white region) and the dc voltage readout amounts $V_1$. The reason for this behavior is a subtle balance between the ac drive and the dc tilt. Exactly at the point were the ac amplitude is strong enough to move the vortex lattice in the ‘hard’ direction (i.e. against the dc tilt) it is also sufficiently strong to move the vortex lattice an additional step in the ‘easy’ direction. Therefore the overall motion of the vortex lattice will be independent on the applied ac drive and equal to 1. This subtle balance is destroyed in favor of ‘easy’ flow if the applied dc current is increased above 100 $\mu$A. Again the dc voltage readout ($1 < V_{dc}/V_1 < 2$) indicates that the vortex lattice moves two steps (maxima in dc voltage) or one step (minima in dc voltage) extra in the positive direction. By increasing the dc tilt even further we could even see a difference of 4 unit cells between the motion in the positive and negative direction.

IX. CONCLUSION

We have demonstrated that the particular geometry of the two-fold symmetric array of magnetic triangles can induce an order-disorder transition as the system switches from the as-grown state to the magnetized state. The disordered phase is characterized by a random distribution of two opposite chirality magnetic buckle states, whereas in the order state all triangles are in a well defined unique magnetic buckle state. The latter phase exhibits unidimensional lines of strong stray fields which represent channels for easy flux motion resulting in a clear anisotropy of the critical current. A quantitative estimation shows that the current needed to induce inter-channels hopping is about 5 times larger than that needed for setting intra-channel motion. It is important to stress that even though we have presented clear evidence of the influence of the ordering of magnetic dipoles induced by the magnetic templating on the vortex dynamics, it is still uncertain whether the ultimate source of disorder in the as-grown state is the random distribution of magnetic poles in the buckle state or the presence of triangles in the magnetic vortex state along one line of connecting triangles. In addition, we showed that the high mobility of the vortex lattice is a consequence of the generation of vortex-antivortex pairs by the magnetic template. This picture was confirmed experimentally by high frequency transport measurements and low frequency ratchet experiments. Due to the self organized commensurability of these channels a pronounced mode locking effect is observed. These features were used to probe and tune the discrete motion of the composite lattice of vortices and antivortices. This combination of techniques represent a powerful tool for understanding the motion of vortices at microscopic scale.

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