"Switching current modulations induced by vortices rearrangement in mesoscopic superconducting loops"

Adam, Sébastien ; Hallet, Xavier ; Piraux, Luc ; Lucot, Damien ; Mailly, Dominique

ABSTRACT

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Switching current modulations induced by vortices rearrangement in mesoscopic superconducting loops

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I. INTRODUCTION

Probing the entrance and the (re)arrangement of vortices in mesoscopic type-II superconducting structures is a longstanding issue, which is still the subject of intensive work.1–5 The reduction of the sample sizes down to the coherence length or the magnetic penetration depth is known to deeply modify the vortex structure, especially due to the competition between the vortex-vortex interaction and the vortex-boundary interaction. Reducing the sample size perpendicular to the applied magnetic field also widens the magnetic-field region over which a sample stays free of vortices.6,7 Indeed, to penetrate into the interior of a superconductor, vortices must overcome the well-known Bean-Livingston surface barrier.8,9 The latter results from the competition between an attractive force toward the surface due to an image antivortex located outside the superconductor, and a repulsive force that comes from the circulating surface currents. This energy barrier vanishes at the so-called critical field for complete vortex expulsion (or more commonly, the superheating field $B_{\text{SH}}$). This characteristic field is inversely proportional to the square of the sample size perpendicular to the applied field$^{10}$ (i.e., $B_{\text{SH}} \propto \Phi_0/W^2$, where $\Phi_0 \equiv \frac{2\pi}{\Phi_1} = 2.07 \times 10^{-15}$ T m$^2$ is the flux quantum and $W$ is the width), so that a mesoscopic superconductor exhibits a complete metastable Meissner state up to a higher magnetic field than the bulk value $B_{1\text{c}}$.

Insofar as the applied magnetic field is higher than the superheating field, a typical field-cooling procedure results in the freezing of the vortices at a temperature that is very close to the critical temperature $T_c$.1,6 The exact location of the vortices depends mainly on the sample size and shape as well as on the number of vortices involved. It is well known that at low fields, the strong repulsive interaction with the circulating currents tends to align the vortices along the center of the strip. This leads to the formation of a one-dimensional chain.11 As the magnetic field is increased, vortices become closer together, leading to an increase of the vortex interaction energy. At some field, the anisotropic lattice becomes unstable and splits into two chains. Similar rearrangements occur at regularly spaced matching fields $B_N$ that correspond to the transition from $N$ to $N + 1$ vortices rows.12 Interestingly, the rearrangement of an unstable vortices pattern into a more stable one goes hand in hand with a minimum of the critical current.13,14 One way to probe the relative stability of the frozen vortices configurations is thus to study their response to a direct transport current.

Recently, Michotte et al.15 have shown on rectangular niobium loops that studying the evolution of the switching current as a function of a perpendicular magnetic field is an efficient tool to probe the fluxoid quantization effect at temperatures far below $T_c$. In particular, they highlighted the impact of the current leads location, either symmetrically or asymmetrically positioned around the loop. This latter point was discussed theoretically by Berdiyorov et al.,16 who mentioned the ability to detect the entry of individual vortex by transport measurements. The field scale $B_{0\Phi} = \frac{\Phi_0}{\pi r_M}$ is given by the area enclosed by the mean radius $r_M$ of the loop ($\Phi_0$ is the flux quantum). This periodicity for the fluxoid quantization is rather different from the one associated with the rearrangement of the vortices configurations, which is directly related to the superheating field.13 These two behaviors can be discriminated unambiguously because the characteristic fields $B_{0\Phi}$ and $B_{\text{SH}}$ depend differently on the sample sizes.

In this paper, the vortices arrangements that occur during the field cooling of thin mesoscopic niobium loops are probed by electrical transport measurements. Periodic modulations of the switching current as a function of a perpendicular magnetic field are observed with a field spacing much higher than the periodicity that governs the fluxoid quantization effect. The local switching current minima are related to unstable vortices
TABLE I. Critical temperature at half the normal state resistance for the loops (Tc,loop) and their contacts (Tc,contact), inner (ri) and outer (ro) radii, branch width (ro-ri), and width of the narrowest contact of the loops L1, L2, and L3. The critical width (i.e., the width of the cross section where the critical current density is first reached) is underlined.

<table>
<thead>
<tr>
<th>Loop</th>
<th>Tc,loop (K)</th>
<th>Tc,contact (K)</th>
<th>ri (nm)</th>
<th>ro (nm)</th>
<th>Branch width (nm)</th>
<th>Contact width (nm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>L1</td>
<td>7.75</td>
<td>8.06</td>
<td>215</td>
<td>365</td>
<td>150</td>
<td>380</td>
</tr>
<tr>
<td>L2</td>
<td>7.75</td>
<td>8.02</td>
<td>420</td>
<td>580</td>
<td>160</td>
<td>300</td>
</tr>
<tr>
<td>L3</td>
<td>7.65</td>
<td>8</td>
<td>255</td>
<td>715</td>
<td>460</td>
<td>400</td>
</tr>
</tbody>
</table>

patterns, whose rearrangement causes a cancellation of the critical voltages (i.e., measured just before the current-driven transition to the fully normal state). In addition, multimodal distributions of the switching current reveal that the few vortices involved in the field-cooling process can freeze into different configurations, even at a single field. Finally, the critical field for complete vortex expulsion is determined from the spacing between the matching fields, in good agreement with recent studies based on scanning probe microscopy.6,7

II. EXPERIMENTAL DETAILS

The superconducting loops were fabricated using the same 20-nm-thick niobium film. The film was obtained by reactive dc magnetron sputtering at room temperature from a pure (99.9%) Nb target at a total Ar pressure of 9.75 mTorr (1000 W dc power). The multicontact structures were obtained by SF6 reactive ion etching using a 40-nm-thick Al2O3 mask patterned with standard e-beam lithography. The loops and the contacts are thus made from the same superconducting material. Typical configurations are shown in Fig. 1. All the experimental results reported in this paper were obtained by using the symmetrical contacts that split up in two (not shown), so that a pseudo-four-contacts arrangement is used to apply the bias and measure the resulting voltage. The electrical transport measurements were performed in a He flow pulse-tube cryocooler with a base temperature of 1.5 K. All the measurement lines were filtered using a combination of RC and LC filters operative up to 40 GHz. The normal-superconductor transition is measured at low direct current and spreads from ~8 K (for the widest parts, i.e., the contacts) down to around 7.6–7.4 K (depending on the loop branch width). The critical temperatures at half the normal state resistance are summarized in Table I for the loops and their contacts. The normal state resistivity determined just above the transition is around 10 μΩ cm.

The Ginzburg-Landau coherence length at zero temperature has been determined from the relation \( \xi(0) = \left[ \Phi_0 / 2\pi B_{c2}(0) \right]^{1/2} \), where \( B_{c2}(0) = 0.69 T_c \left| \frac{dB_c}{dT} \right|_{T=T_c} \) is the reduced upper critical field at zero temperature.17 The temperature dependence of \( B_{c2} \) has been monitored for each loop in a perpendicular magnetic field up to 3 T. The expected linear dependence at \( T \approx T_c \) is verified for all the samples18 and one obtains that \( \xi(0) \approx 9.5 \) nm. Furthermore, the magnetic penetration depth at zero temperature is expressed in the dirty limit by the Bardeen-Cooper-Schrieffer expression \( \lambda(0) = \sqrt{\frac{\hbar}{\pi \rho_n \mu_0 \Delta(0)}} \), where \( \hbar \) is the reduced Planck constant, \( \rho_n \) is the normal-state resistivity, \( \mu_0 \) is the vacuum permeability, and \( \Delta(0) \) is the energy gap at zero temperature. Niobium is in the intermediate-coupling limit with the superconducting energy gap given by \( 2\Delta(0) \approx 3.6 k_B T_c \) (\( k_B \) is the Boltzmann constant).19,20 For our samples, the energy gap at zero temperature is around 1.2 meV so that the magnetic penetration depth at zero temperature amounts to ~120 nm, in agreement with the results of Gubin et al.21 In addition, in accordance with the

![FIG. 1. Scanning electron microscope images of 20-nm-thick niobium loops. The surrounded areas highlight the weak regions, i.e., the regions where the critical current density is first reached under a direct applied current.](image_url)
SWITCHING CURRENT MODULATIONS INDUCED BY ...

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FIG. 2. (a) Evolution of the switching current of the loop L1 as a function of a perpendicular magnetic field. The bath temperature (3 K) lies far below the critical one. (b) Enlargement of the rectangular frame of (a). The field values \( B_N \) that correspond to local minima of the switching current are numbered for \( N = 1, 2, \ldots \). The symbols [A] and [B] refer to two particular magnetic fields studied in more details in Fig. 3.

FIG. 3. Numerous switching current measurements for the loop L1 obtained at fixed magnetic fields ([A], 193 mT; [B], 265 mT). A total of 500 measurements have been performed, and only the first 150 are shown. Three well-defined values are observed at each field. The narrow distributions are depicted on the right side.

The loop L1 [Fig. 1(a)] has a contact width larger than twice the branch width so that the weak region (i.e., the region where the critical current density is first reached) is located in one of the two branches. In contrast, the loops L2 and L3 [illustrated for L2 in Fig. 1(b)] have a contact width smaller than twice the branch so that the critical area is close to the contact. The critical widths \( W \) to take into account are underlined in Table I and amount to 150 nm (for L1), 300 nm (for L2), and 400 nm (for L3).

III. RESULTS

The overall evolution of the switching current \( I^*_c \) under a perpendicular magnetic field is given in Fig. 2(a) for the loop L1. At each magnetic-field step, the sample is brought to the normal state by the applied direct current so that the critical electronic temperature is exceeded under the influence of the Joule heating. The sample is then field-cooled when the current source is switched off, although this is not a temperature-driven field cooling as the bath temperature is maintained constant during the whole measurement. The magnetic-field increment is 100 \( \mu \)T and measurements have been performed in both the positive and negative orientations. Figure 2(a) is representative of the low-temperature behavior for all the samples (here \( T = 3 \) K, i.e., well below the critical temperature of the loop). Figure 2(b) depicts in more detail the wide variety of switching current modulations that appear in the rectangular frame of Fig. 2(a). A relevant feature is that the local minima are regularly spaced, with a period that amounts to around 65 mT (the magnetic-field values \( B_N \) associated with the local switching current minima are numbered for \( N = 1, 2, \ldots \)). Here, the period of interest is much larger than that observed in the low-field part of the curve, where a sawtooth-shaped evolution of the switching current arises with a periodicity of 7.8 mT. This latter value is in excellent agreement with the periodicity of the fluxoid quantization [i.e., \( B_{\Phi 0} = \frac{\Phi_0}{\pi \lambda(0)^2} \approx 7.83 \text{ mT} \)], which is calculated with the mean radius of the loop.
FIG. 4. (a) Evolution of the switching current of the loop L2 as a function of a perpendicular magnetic field. The bath temperature (4 K) lies far below the critical one. (b) Enlargement of the rectangular frame of (a). The field values $B_N$ that correspond to local minima of the switching current are numbered for $N = 3, 4, \ldots$

L1. The modulations observed in the low-field region are very similar to the one recently observed on nonoxidized niobium nanosquids. It is also noteworthy to measure very different values of the switching current at a given magnetic field (see, for example, the fields denoted $A$ and $B$ in Fig. 2(b), where $I_s$ spreads from 280 to 360 $\mu$A and from 220 to 270 $\mu$A, respectively). The measurements have been reproduced 500 times at these particular fields (see Fig. 3, in which only the first 150 are shown). The three critical current levels pointed out at each field nicely match the overall curve. The spreading is very narrow, and the lowest values appear to be more frequently probed. Multimodal distributions at fixed magnetic field reveal that the sample is randomly stabilized in different states during the current-driven field cooling that follows each critical current detection (i.e., a larger switching current is related to a better stabilized starting state).

The behavior depicted in Fig. 2 for the loop L1 has been similarly observed on the loops L2 and L3. This is highlighted for the sample L2 in Fig. 4 with the overall evolution of the switching current [Fig. 4(a)] and an enlargement of the rectangular frame in the negative field orientation [Fig. 4(b)]. The bath temperature is 4 K. Once again, the regularly spaced $I_s$ minima can be numbered. This is done for $N = 3, 4, \ldots$ as the two first minima are located in the central peak, which is not enlarged in Fig. 4(b). The field period amounts to $\sim 33$ mT, around half the value determined for the loop L1. Compared to the loops L1 and L2, the loop L3 is characterized by broader distributions of the switching current (see Fig. 5). The local minima (numbered from $N = 1$ to 5) remain, however, clearly defined with a smaller field period of $\sim 18$ mT.

Contrary to the loop L1, for which the fluxoid quantization manifests itself in the low-field part of the curves, no similar feature has been observed for the loops L2 and L3. This is consistent with the position of the critical width, which is located either in a branch of the loop (for L1) or near the contacts (for L2 and L3). Indeed, the circulating currents that ensure the quantization of the fluxoid flow in the branches of the loops. They subtract from the applied transport current in one branch and add in the other one where the switching current is first reached. The situation alternates periodically from one branch to the other with a total field period given by $B_{B_0}$. The sawtooth evolution of the switching current (with respect to the applied field) thus directly indicates the presence of the circulating currents in the branches of the loops (see, e.g., the work of Michotte et al. on the evolution of the switching current of a niobium superconducting loop far below the critical temperature). For the loops L2 and L3, the presence of the circulating currents is not probed in the magnetic-field evolution of the switching current because they do not impact the net current that flows through the contacts. If the fluxoid quantization effect would be probed, the theoretical field period would be noticeably less than the ones involved in Figs. 4 and 5. Indeed, $B_{B_0}$ would amount to 2.6 mT for L2 and to 2.9 mT for L3, while the observed field spacing between the switching current minima is 33 and 18 mT for the loops L2 and L3, respectively.

FIG. 5. Evolution of the switching current of the loop L3 as a function of a perpendicular magnetic field. The bath temperature is 3 K. The field values $B_N$ that correspond to local minima of the switching current are numbered for $N = 1$ up to $N = 5$. 

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The evolution of the switching voltage $V^*$ with respect to the magnetic field is also instructive. This is illustrated (together with the switching current $I^*$) in Figs. 6(a) and 6(b) for the loop L1 and L2, respectively. The switching voltage, which is measured just before the transition to the fully normal state, takes either zero or nonzero values depending on the magnetic-field range. The nonzero values are most often probed when the switching current is decreasing, and $V^*$ generally drops at the local $I^*_c$ minima. The ability to relate both $I^*_c$ and $V^*$ is the principal advantage of using a switching criterion rather than a criterion based on the appearance of resistance (or voltage). Furthermore, it has to be mentioned that the evolutions of $I^*_c$ and $V^*$ are reproduced nicely, either in an increasing or in a decreasing magnetic field. As was previously emphasized, the electrons are briefly heated above $T_c$ each time $I^*_c$ is reached, so that the sample undergoes a kind of thermal reset after each switching current detection. Therefore, the results are expected to be independent of the magnetic-field sweeping direction.

IV. DISCUSSION

The response of a type-II superconductor under a direct transport current depends mainly on the stability of the vortices pattern that freezes during the field-cooling procedure. In the mixed state, the vortices are forced to move when the Lorentz driving force, which is proportional to the applied transport current density, exceeds a certain pinning threshold. If the threshold is small (i.e., the vortices configuration is not very stable), the vortices move at relatively low applied currents. The dissipation induced by their motion destroys the perfect superconducting state: a nonzero voltage appears at the extremities of the sample. A further increase of the applied current is required in order to entail the transition to the normal state. On the other hand, a large transport current must be applied when the pinning threshold is enhanced (i.e., for a more stable vortices configuration). In this case, the dissipation that results from the motion of a single vortex may be sufficient to induce the transition to the normal state, with a zero voltage measured just before the transition.

The experimental results presented in Sec. III reflect the above-described stability of the vortices arrangements that form in the arms and in the contacts of the thin mesoscopic niobium loops. The modulations of the switching current highlight that more or less stabilized initial states are formed during the field cooling that precedes each measurement. This point is corroborated by the correspondence between $I^*_c$ and $V^*$ (see Fig. 6). Indeed, nonzero values of $V^*$ are observed when the switching current decreases. In these field intervals, less stable vortices configurations develop and the vortices are easily displaced by the direct current. On the contrary, zero switching voltages mean that more stable configurations arise. As was already mentioned, the critical current is the lowest when the unstable vortices system rearranges. In this sense, the sharp variations of $V^*$ that occur in the vicinity of the $I^*_c$ minima are of particular interest because they reinforce the occurrence of a transition from an unstable vortices configuration ($V^*_c \neq 0$) into a more stable one ($V^*_c = 0$) at the particular magnetic fields $B_N$. The dependence of these latter fields with respect to the number $N$ of rows is summarized for the three samples in Fig. 7 (the data are shown for both the positive and the negative orientations of
the applied field). The different slopes result from the varying spacing between the fields \( B_S \), which amounts to 65, 33, and 18 mT for loops L1, L2, and L3, respectively. The evolution of \( B_N \) with the number of vortices lines has been discussed in detail by Ziese et al.\(^{13}\) In the limit \( \xi/\gamma W \ll 1 \) (with \( \gamma \) an anisotropy factor that is close to 1 here and \( W \) the sample size perpendicular to the applied field), the fields \( B_N \) depend linearly on \( N \) according to Eq. (1),

\[
B_N = B_{SH}(0.19 + 0.71N),
\]  

where \( B_{SH} \) is the superheating field, i.e., the maximum field for complete vortex expulsion without an applied transport current. One sees that the formation of the first vortices line is made easier under the action of a transport current (i.e., the critical field for complete vortex expulsion). Two theoretical models are based on the Gibbs free-energy profile of an isolated vortex located in a strip, while a third one deals with the formation of a vortex-antivortex pair. The first (metastable) solution argues that a vortex can penetrate into a strip insofar as the magnetic-field induction is large enough to induce a minimum of the Gibbs free energy at the center of the strip. This occurs for fields \( B > B_0 \), with \( B_0 \) given by Eq. (2),

\[
B_0 = \frac{\pi \Phi_0}{4W^2},
\]  

where \( W \) is the sample width. A further increase of the induction is required in order to induce an absolutely stable solution. This arises at a magnetic field \( B_s \) given by Eq. (3),

\[
B_s = \frac{2}{\pi} \frac{\Phi_0}{W^2} \ln \left( \frac{\alpha W}{\xi} \right),
\]

at which the Gibbs free energy at the center of the strip presents an absolute minimum (\( \alpha \) is a constant that takes the value 2/\( \pi \) or 1/4 depending on the assumption on the vortex core size). Recently, Kuit et al.\(^{1} \) proposed an intermediate model that takes into account the vortex-antivortex pair production energy. They show that the relevant critical induction is determined when the vortex generation rate is exactly balanced by the rate of escape, which occurs at a magnetic field \( B_K \) given by the following expression:

\[
B_K = 1.65 \frac{\Phi_0}{W^2}.
\]  

A common feature of these models is that the superheating field is roughly proportional to \( \Phi_0/W^2 \). The \( B_{SH} \) values extracted here from electrical transport measurements are fully consistent with this predicted dependence on the sample width. This is illustrated in Fig. 8(a). The superheating field of loops L1, L2, and L3 are indicated with the full symbols, while the critical inductions experimentally determined by Stan et al.\(^{6} \) with a scanning Hall-probe microscope are indicated with the empty symbols. These values correspond to the maximum field at which the vortices are completely expelled by thin niobium strips whose width is in the \( \mu \)m range. One notices that the \( B_{SH} \) values extracted here follow nicely the general trend given (i) by the wider niobium strips and (ii) by the theoretical models. Indeed, the solid curve in Fig. 8(a) corresponds to the relation \( k \Phi_0/W^2 \) with the proportionality constant \( k = 1.8 \).
Figure 8(a) compares the experimental data with the three theoretical models. The proportionality term does not depend on the sample width for the $B_0$ (metastable) and the $B_K$ models, so that they are both depicted by straight lines in this log-log plot (see the dotted curve for $B_0$ and the solid curve for $B_K$). Besides, the stable solution $B_s$ depends on the coherence length at the temperature $T_f$, at which the vortices become pinned during the field cooling. Many authors have indicated that the temperature $T_f$ is very close to the critical one in a temperature-driven field cooling ($T_f \approx 0.99 T_c$). However, whether this statement remains valid here is not so clear. Indeed, our mesoscopic superconductors are field-cooled by switching off the applied current at a fixed bath temperature so that the (unknown) temperature at which the vortices become pinned may be reduced or may even be close to the bath one. For this reason, the stable solution $B_s$ (with $\alpha = 2\pi$) is depicted for different $T_f/T_c$ ratios, namely 0.99, 0.9, and 0.4 (see the dashed curves, labeled A, B, and C). This of course corresponds to different values of the coherence length $\xi(T_f)$ that are determined using the temperature dependence of the coherence length in the dirty limit, \[ \xi(T) = \left( \frac{T_0}{T - T_0} \right)^{0.5} \] with $\xi(0) = 0.855 (\xi_0)^{0.5}$. One obtains that $\xi(T_f) \approx 95$, 30, and 12 nm, respectively. Apart from the curve that corresponds to a pinning temperature very close to $T_c$, the stable solution $B_s$ satisfactorily follows the experimental data. This is also the case for the $B_K$ model, so that it appears to be very difficult to discriminate between these two models. Moreover, the metastable solution $B_0$ seems to underestimate slightly the results in the whole studied width range. A definitive conclusion concerning the theoretical model that is better suited to describe our experimental data would require us to know with precision the size of the vortices cores [$\sim \xi(T_f)$] at the temperature at which the vortices positions are frozen. This information is, however, not accessible in our measurements. In any case, the overall agreement between the theory and the experimental data highlights the fact that studying the matching fields at which the vortices lattice rearranges itself is an interesting tool to determine the superheating field of thin mesoscopic superconducting structures with a width of a few hundreds of nanometers (and probably less).

V. CONCLUSION

This paper highlights the fact that electrical transport measurements are efficient tools to probe the vortices arrangement in superconducting mesoscopic niobium structures. The evolution of both the switching current and voltage (i.e., measured just before the current-driven transition to the fully normal state) was studied as a function of a perpendicular magnetic field. The modulations in the switching current confirm that the vortices rearrange themselves at regularly spaced matching fields. Here we highlight that this rearrangement from an unstable vortices pattern into a more stable one also causes a drop of the switching voltage. The latter takes either zero values (for stable vortices configurations) or nonzero values (for less stable vortices configurations) depending on the stability of the vortices patterns under a direct transport current. The observed multimodal distributions of the switching current obtained at a constant field show that the few vortices involved in the measurements can freeze into different configurations during the current-driven field cooling that precedes each measurement.

Furthermore, the spacing between the matching fields allows one to extract the critical field for complete vortex expulsion of the samples, i.e., the so-called superheating field. The values are compared with several critical induction models, which describe the dependence of the superheating field with the sample width. The values extracted here from electrical transport measurements are fully consistent with the predicted evolution $B_{\text{SH}} \propto \Phi_0/W^2$. Both the stable solution and the model based on the vortex-antivortex pair production energy appear well-suited to explain the results. However, a definitive discrimination is difficult due to the unknown temperature at which the vortices become pinned in the field-cooled measurement. In addition, our data are in nice agreement with recent values obtained by scanning probe microscopy on wider niobium strips. Here, the main advantage is that the superheating field is determined for much narrower samples, with no limitation due to the experimental setup resolution.

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23. The curve asymmetry with respect to positive and negative field orientations is believed to reflect the asymmetry of the sample configuration. Indeed the third contact is thought to locally modify the transport current density, although it is not used in the measurement.