Nonlocal vortex motion in mesoscopic amorphous Nb$_{0.7}$Ge$_{0.3}$ structures

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We study nonlocal vortex transport in mesoscopic amorphous Nb$_{0.7}$Ge$_{0.3}$ samples. A dc current $I$ is passed through a wire connected via a perpendicular channel, of a length $L = 2 - 5 \mu$m, with a pair of voltage probes where a nonlocal response $V_{nl} \propto I$ is measured. The maximum of $R_{nl} = V_{nl}/I$ for a given temperature occurs at an $L$-independent magnetic field and is proportional to $1/L$. The results are interpreted in terms of the dissipative vortex motion along the channel driven by a remote current, and can be understood in terms of a simple model.

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In a pioneering work Giaver measured a magnetic-flux-transformer effect in type II superconductors. He applied a magnetic field $B$ perpendicularly to a sample comprising two superconducting sheets separated by a thin insulator, passed a current $I$ through one of the superconductors and measured a voltage developed over the other one - where no current was flowing. The induced voltage was a consequence of an electromagnetic coupling of vortices in the two layers. In their recent experiment Grigorieva et al. demonstrated a complementary flux-transformer phenomenon associated with vortices. They produced mesoscopic amorphous MoGe structures of a double-cross shape, consisting of two parallel wires connected at a right angle by a channel of a width $w = 0.07 - 2 \mu$m and a length $L = 0.5 - 12 \mu$m. In a perpendicular $B$ and with $I$ through one of the parallel wires a nonlocal voltage $V_{nl}$ appeared over the second, current-free wire. This novel, transversal flux-transformer effect originated in the in-plane vortex-vortex repulsion which conveyed the driving force from the current-carrying wire to the vortices in the channel. The effect disappeared not only for $L$ exceeding 6-7 $\mu$m but also for $w$ larger than $\sim 0.5 - 1 \mu$m. When $w$ was sufficiently small the force on the vortices in the channel was transferred over many intervortex distances, and, moreover, $V_{nl}$ was proportional to $I$. The efficiency of transversal flux-transformer effect can be quantified by a nonlocal resistance $R_{nl} = V_{nl}/I$.

In the experiment of Grigorieva et al. the local mixed-state dissipation was characterised on separate mm-sized films, whereas $V_{nl}$ was measured by a low-frequency ac method during $B$ sweeps at constant temperatures $T$. An ac method was used because $V_{nl}$ was in nV range, i.e. $R_{nl} < 5 \text{ m}\Omega$, thus being too small for a dc detection. In our work we focused on dc probing of transversal flux-transformer effect and measuring $V_{nl}$ and the local voltage $V_i$ on a same sample, which was possible in multi-terminal amorphous (a-)Nb$_{0.7}$Ge$_{0.3}$ structures of the geometry shown in the inset to Fig. 1. The weak pinning, characteristic of the a-Nb$_{0.7}$Ge$_{0.3}$ material used, resulted in a dc-measurable $V_{nl}$ and $R_{nl} \sim 1 \Omega$ even at very low temperatures. The measured nonlocal resistance was hence two orders of magnitude larger than in Ref. 2.

We measured the transversal flux-transformer effect in samples of different length by isothermal sweeps of $B$ at different applied $I$ and for $0.15 T_c \lesssim T \lesssim 0.95 T_c$, where $T_c$ is the superconducting transition temperature. $V_{nl}$ depends linearly on $I$ in the range $I = 0.1 - 1 \mu$A. With increasing $B$, $R_{nl}(B)$ first acquires a nonzero value at a certain magnetic field $B_d$, then has a maximum at some $B_p$ and gradually vanishes close to the upper critical magnetic field $B_c$. The main representatives of $R_{nl}$, i.e. $B_p$ and $R_p = R_{nl}(B_p)$, behave differently with respect to the channel length. $B_p$ is independent of $L$ whereas $R_p \propto 1/L$, suggesting a vortex velocity $u_{nl} \propto 1/L$ at the nonlocal voltage probes. This we relate to the total frictional force on the vortices in the channel being proportional to $L$.

We investigated two structures of the type shown in the inset to Fig. 1, where we also assign numbers to the leads and define the coordinate system (with the unit vectors $\hat{x}$, $\hat{y}$ and $\hat{z}$). As in our previous studies, the samples were produced by combining electron-beam lithography with magnetron sputtering but the film thickness was increased from 20 nm to $d = 60$ nm in order to safely avoid inhomogeneities detected in samples of a cross section smaller than $\sim 5000$ nm$^2$. This way the pinning was enhanced, as we inferred from the local dissipation, but still remained weak enough to permit a nonlocal vortex motion over several microns. Contacts 0 and 1 are used for applying a dc $I = I_{01}$ through the horizontal wire, exerting a local force on vortices in the $y$ direction (since $B = -B\hat{z}$). Having a velocity $\mathbf{u} = u\hat{y}$, vortices induce an electric field $\mathbf{E} = E\hat{x}$. Combinations of contacts $i, j = 6 - 9$ are used for measuring the local voltage drop $V_i$. The local vortex pressure is transferred from the horizontal wire along the zero-$I$ channels (the length $L$ is indicated for the left channel) toward the crosses contacted by leads 2,3 and 4,5, where $V_{nl}$ is measured. For the first sample the channel lengths were 2 and 3 $\mu$m, and for the second 3 and 5 $\mu$m, so we covered the range $L = 2 - 5 \mu$m and had two $L = 3 \mu$m samples
for consistency check. All wires of a single sample had the same width, \( w = 275 \) nm for the sample with the longer channels and \( w = 250 \) nm for the other one. This small difference did not affect the results presented henceforth. Since dc measurements are invariably burdened by sub-\( \mu \)V parasitic signals, in order to determine \( V_{\text{nl}}(I) \) for a given measurement with leads \( i,j \) we recorded both \( V_{\text{nl},ij}(I) \) and \( V_{\text{nl},ij}(I = 0) \) taken at the same \( B \)-sweep rate and direction. \( V_{\text{nl},ij}(I = 0) \) was different for different pairs of leads and depended very weakly on the \( B \)-sweep rate. The background-free nonlocal voltage was extracted as \( V_{\text{nl}}(I) = V_{\text{nl},ij}(I) - V_{\text{nl},ij}(I = 0) \). By this procedure we found a very regular behavior of \( V_{\text{nl}}(I) \), which, in particular, for the two \( L = 3 \) \( \mu \)m samples agreed within the error bars. From the local voltage response between contacts (8,9) we characterised the samples in the same way as in our previous work obtaining \( T_c = 2.95 \) K, the normal-state resistivity \( \rho_0 = 3.7 \mu \Omega \)m, \(- (dR_{\text{nl}}/dT)_T = 2.17 \) T/K, and the Ginzburg-Landau parameters \( \kappa = 100, \xi(0) = 7.15 \) nm and \( \lambda(0) = 1.18 \) \( \mu \)m.

Results typical of \( V_{\text{nl}} \) are shown in Figs. 1 and 2. The data in Fig. 1 were obtained for the \( L = 2 \) \( \mu \)m channel at \( T = 1.4 \) K (\( B_{\text{c2}} = 2.95 \) T) and with \( I = I_{\text{d1}} = 300, 600, 800 \) nA. The lines (left-hand scale) correspond to the local response \( V_{\text{sq}}/I_{\text{d1}} \), and the symbols (right-hand scale) to \( V_{\text{nl}}(I)/I = [V_{\text{sq}}(I_{\text{d1}}) - V_{\text{sq}}(0)]/I_{\text{d1}} \). As expected from previous investigations of \( V(I) \) of \( \alpha-\text{NbSe}_2 \) samples, the current density \( J \sim 20-50 \) MA/m\(^2\) the local response at low \( T \) depends on \( I \). Noteworthy, in contrast to \( d = 20 \) nm samples with a weaker pinning, the onset of \( V_{\text{nl}} \) at a certain \( B = B_{\text{d}} \) is essentially independent of \( I \). \( B_{\text{d}} \) nearly coincides with \( B_d \) for \( L = 2 \) \( \mu \)m, while \( B_d \) for \( L = 3.5 \) \( \mu \)m is higher for \( \sim 5 \) % (high \( T \)) to \( \sim 20 \) % (low \( T \)) and mutually indistinguishable. Contrary to \( V_1(I) \), \( V_{\text{nl}}(I) \) is for the given range of \( I \) linear for all \( T \).

In Fig. 2 we show how \( R_{\text{nl}} = V_{\text{nl}}/I \) changes with temperature. The results refer to an \( L = 3 \) \( \mu \)m channel, \( I = 600 \) nA and \( T = 0.8, 1.6, 2.6 \) K. As in Fig. 1, \( R_{\text{nl}} = 0 \) up to \( B = B_{d2} \) after which it displays a relatively broad peak around \( B = B_p \), which defines \( R_p = R_{\text{nl}}(B_p) \). Close to the \( B_{d2}, R_{\text{nl}} \) drops to zero again. While the nonlocal resistance in Ref.\(^2\) decreased with decreasing \( T \), vanishing at \( T/T_c \sim 0.6 \), in our experiment we find a nonmonotonic variation of \( R_{\text{nl}}(T) \). \( R_{\text{nl}} \) is finite even at the lowest measurement temperature \( (T = 0.4 \) K, \( T/T_c \approx 0.14 \)), decreases at low \( T \) and increases close to \( T_c \) with increasing \( T \) (see Fig. 3(b) later). The \( R_{\text{nl}}(B) \) traces are nearly symmetric around \( B = 0 \) but, especially at low \( T \) and/or high \( I \), \( R_p \) for \( B < 0 \) and \( B > 0 \) may differ to some extent, as seen, e.g., in Fig. 1 for \( I = 800 \) nA (1.4 K) and in Fig. 2 for \( T = 0.8 \) K (600 nA). These differences do not necessarily appear only because of variations in measurement conditions during the long \( B \) sweeps (e.g., slightly different bath temperature), they may also originate in the Nernst effect due to a heating in the current-carrying wire. However, the Nernst effect changes sign by reversal of \( B \) and can be canceled by the averaging \( R_{\text{nl}}(B) = [R_{\text{nl}}(+B) + R_{\text{nl}}(-B)]/2 \). Moreover, for our samples \( R_{\text{nl}}(+I) = R_{\text{nl}}(-I) \), which rules out a rectifying effect proposed recently.

Essential information on the nonlocal vortex motion is contained in the characteristic quantities \( B_p \) and \( R_p \). In Fig. 3(a) we plot \( B_p(T) \) for all three channel lengths, together with \( B_{d2}(T) \), \( B_{dl}(T) \) and \( B_p(T) \) for the \( L = 2 \) \( \mu \)m channel. \( B_{dl} \) and \( B_p \) are relatively low and become immeasurably small above 2.6 K (this applies to \( B_p \) for \( L = 3.5 \) \( \mu \)m as well). It can be seen that \( B_p \) is independent of \( L \). On the other hand, \( R_p \) does depend on \( L \), which is shown in Fig. 3(b). \( R_p \) vs \( T \), plotted by the symbols, exhibits a reasonably-well-defined scaling behaviour implying \( R_p \propto 1/L \).

We shall now discuss a simple model that leads to a
reasonable explanation of the presented experimental results. This model assumes equilibrium between the driving force exerted by $I$ on the vortices in the lower cross and the frictional damping force on the vortices in the channel, as well as the presence of surface barriers enforcing a dominating $y$ component of $u$ inside the channel. In a first step we neglect the pinning and assume that the frictional force is linear in $u$ with a velocity-independent friction coefficient $\eta$. The effects of pinning are addressed later.

At a vortex density $n_0 = B/\phi_0$, where $\phi_0$ is the magnetic-flux quantum, in total $n_0 w^2$ vortices in the lower cross each experiencing a force $J\phi_0 d$, apply a pressure $p = n_0 \phi_0 I/d$ on the vortices in the channel. The corresponding pushing force (per unit vortex length) $pw$ is balanced by the force required to move $n_0 w L v$ vortices along the channel against the frictional damping $\eta u n_l$ per vortex. This gives $u n_l = \phi_0 I / \eta L d \propto 1/L$. Using $V_{nl} = wB u n_l$ we find

$$R_{nl} = \frac{V_{nl}}{I} = \frac{\phi_0 B w}{L \eta d} \quad .$$

The above result holds if the surface barriers are strong enough to confine the vortex motion within the channel. As shown experimentally in Ref. 2, this is not satisfied for large $w$. In any case, surface barriers are essential to preserve the uniaxial character of the nonlocal vortex motion. An important source of the surface barriers are the Meissner currents $J_M$ flowing along the channel edges and providing an inward-pointing force $F_{in} \propto J_M$. The surface barriers weaken by approaching $B_{c2}$ irrespectively of their exact origin, which my explain the fact that $R_{nl}(B)$ does not increase all the way up to $B = B_{c2}$ before dropping to zero in the normal state. This restricts the range of the applicability of Eq. (1) to $B$ not too close to $B_{c2}$. On the other hand, Eq. (1) correctly reproduces $R_{nl} \propto 1/L$ and, as we show below, accounts for $R_{nl}(I)$ quantitatively in conditions of insignificant pinning. For $T \geq 2.7$ K the $V_I(I)$ curves are linear beyond any doubt, implying a negligible pinning and, moreover, $\eta \approx \eta_I$ of the viscous drag in pure flux flow. There are two possible dissipation mechanisms that determine the flux-flow viscosity $\eta_I$. One is related to Joule heating of normal electrons in vortex cores by the $E$ therein, as described by the Bardeen-Stephen model giving $\eta_I = \eta_{BS} = \phi_0 B_{c2}/\rho_n l_{ab}$. The other approach, proposed by Tinkham, attributes the dissipation to a loss of the superconducting Gibbs free-energy density $G_s(B, T)$ as vortices move and cause depairing and recombination of Cooper pairs. This process is affected by the time $\tau = (\hbar/\Delta_0)[1 + (T/T_c)^2]/[1 - (T/T_c)^2]$ of establishing a superconducting state and results in $\eta_I = \eta_{BS} = 2\pi \tau G_s$ ($\Delta_0 \approx 1.76 k_B T_c$ is the superconducting gap at $T = 0$, and $h$ the Planck constant). For $B$ not too close to $B_{c2}$, i.e. in the regime of the applicability of Eq. (1), $G_s$ can be approximated by the superconducting condensation energy $U_s(T) \approx B_s^2(T)/4\kappa^2 \mu_0$, where $\mu_0 = 4\pi/10^7$ H/m. Since we know $B_{c2}(T)$ we can determine both $\eta_{BS}$ and $\eta_I$ for our samples, which permits a quantitative comparison of the experiment and the model.

At $T \gtrsim 2.7$ K the measured $B_p \approx 60$ mT $\ll B_{c2}$ is fairly constant, so the temperature dependence of the corresponding $R_p$ is dominated by that of $\eta$. By calculating $\eta_{BS}$ and $\eta_I$ without any adjustable parameter and using Eq. (1) we obtain $LR_p$ shown in Fig. 3(b) by the solid lines, as indicated. $R_p(T)$ is well reproduced in both cases. In particular, the steep increase of $R_p$ near $T_c$ can now be traced back to the rapid decay of $\eta_I$ close to $T_c$.

So far we have not included any effect of pinning into our analysis, which is justified only in a narrow $T$ range close to $T_c$. At lower temperatures the use of $\eta_I$ and the measured values of $B_p$ in Eq. (1) results in a rapid increase of $R_{nl}$ far above the observed values. In addition, nonlinearities in $V_I(I)$ are observed below 2.7 K, indicating that the pinning is no longer marginal. However, the fact that the linearity of $V_{nl}(I)$ and $R_p \propto 1/L$ are preserved in this regime brings up a possibility of extending our model to lower temperatures. This is achievable if even in the presence of pinning $p$ remains proportional to $I$ and $\eta$ independent of $u$. Below we argue that these properties are consistent with a plausible picture of the distribution of the pinning force in our samples.

We recall that the onset of $V_I$ at $B_{ab}$ does not depend on $I$, which suggests a critical vortex density for triggering the dissipation. This implies that for $B < B_{ab}$ the flux is first trapped in the most-strong-pinning regions of the channels.
When these are saturated at $B = B_{dl}$ vortices enter the lower-pinning regions in between - where they can move more easily. Only after this saturation the dissipation starts by the vortex motion along "vortex rivers" between the strong-pinning sites/areas. In the current-carrying wire vortices in the "rivers" shear plastically with the edges, thus enhancing the driving force above $J\phi_0$ in depinning the vortices. This process may result in a nonlinear $V(I)$ such as that observed in Fig. 1 but cannot propagate significantly into the channel, where the external driving current decays exponentially. The pushing force is thus conveyed along the "rivers" set by the equilibrium pinning properties, and the fraction of the contributing vortices is $f \sim (B - B_{dl})/B$. The vortex pressure is therefore reduced by the same factor $f$.

The frictional damping of the vortex motion in "vortex rivers" has been investigated in samples with artificial easy-flow channels embedded in a strong-pinning medium. Because of the random pinning landscape in our samples no commensurability effects between the moving and immobile vortex regions are expected and the vortex velocity in the "rivers" should respond linearly to the applied force. This may explain the observed linearity of $V_{nl}(I)$, while the magnitude of $R_{nl}$ is reduced by the ratio $f$ that accounts for the number of vortices in the weak-pinning regions. Although we lack a manageable model for calculating $\eta$ including the effects of pinning, the increase of $R_p$ as $T$ is lowered can be explained at least qualitatively as a consequence of $B_p(T)$ growing faster than $\eta(T)$.

To conclude, we have investigated a transversal flux-transformer effect, which manifests itself in a nonlocal flow of vortices in a narrow superconducting channel driven by a remote dc current. In our low-pinning $\alpha$-Nb$_{0.7}$Ge$_{0.3}$ the effect appears in more than half of the superconducting phase diagram, i.e. everywhere where the vortices can be moved easily enough to induce dissipation at very low currents. The effect is two orders of magnitude larger than in previous studies and we observe a nonmonotonic variation of the maximal nonlocal resistance with temperature, which can be explained by the interplay of vortex density and vortex viscosity. Close to $T_c$ the data are in quantitative agreement with a simple model which is based on the assumption that the vortices behave like a weakly compressible fluid confined to the superconducting channel by surface barriers.

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