Collective Josephson vortex dynamics in a finite number of intrinsic Josephson junctions

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We report the experimental confirmation of the collective transverse plasma modes excited by the Josephson vortex lattice in stacks of intrinsic Josephson junctions in Bi2Sr2CaCu2O8+x single crystals. The excitation was confirmed by analyzing the temperature (T) and magnetic field (H) dependencies of the multiple sub-branches in the Josephson-vortex-flow region of the current-voltage characteristics of the system. In the near-static Josephson vortex state for a low tunneling bias current, pronounced magnetoresistance oscillations were observed, which represented a triangular-lattice vortex configuration along the c axis. In the dynamic vortex state in a sufficiently high magnetic field and for a high bias current, splitting of a single Josephson vortex-flow branch into multiple sub-branches was observed. Detailed examination of the sub-branches for varying H field reveals that sub-branches represent the different modes of the Josephson-vortex lattice along the c axis, with varied configuration from a triangular to a rectangular lattices. These multiple sub-branches merge to a single curve at a characteristic temperature, above which no dynamical structural transitions of the Josephson vortex lattice is expected.

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I. INTRODUCTION

It has been known that stacked intrinsic Josephson junctions (IJJs) form in naturally grown Bi2Sr2CaCu2O8+x (Bi-2212) single crystals. The collective Josephson plasma oscillation, manifested by the electromagnetic resonant absorption in such Josephson-coupled layered superconductors, has provided the key to understanding the superconducting state and the vortex phases forming in the materials. Two kinds of collective Josephson plasma modes, longitudinal and transverse, exist in such a system of an infinite number of stacked IJJs. The longitudinal plasma mode oscillation, excited in an externally applied c-axis-oscillating electric field, propagates along the c axis of a stack of IJJs with the inter-layer phase difference of a junction being uniform in a planar direction. The resonant absorption of externally applied microwaves by the excitation of the longitudinal modes enables one to directly obtain the Josephson plasma frequency, ωp. This, in turn, provides information on the interlayer Josephson coupling strength and the c-axis superfluid density in Bi-2212. By contrast, the transverse plasma modes can be excited by the oscillating magnetic fields in the ab plane. The resulting induced current along the c axis generates the oscillating electric field that resonates with the transverse Josephson plasma oscillation.

On the other hand, a system with a finite number (N) of stacked IJJs, described by linearized coupled sine-Gordon differential equations with the inductive interjunction coupling, exhibits collective transverse plasma (CTP) eigen modes, with the eigen frequencies lying between those of the longitudinal and the transverse modes found for a system of an infinite number of stacked junctions. The corresponding dispersion relation is expressed as

$$\omega_n(k) = \omega_p \sqrt{1 - \frac{\lambda_c^2 k^2}{1 + (2\lambda_{ab}/sD)(1 - \cos(\frac{m}{N+1}))}}. \quad (1)$$

Here, λab (λc) is the c-axis (ab-plane) London penetration depth, N the number of the junctions in a stack, s=0.3 nm (D=1.2 nm) the thickness of a superconducting (insulating) layer, k(=2πm/L; L=the length of a stack; m=the number of the vortices present in a junction) the wave number of the plasma oscillation. The mode index n runs from 1 to N.

The CTP modes can be excited by the moving Josephson vortex lattice forming in high magnetic fields, which are alternative solutions of the coupled sine-Gordon equations. The temporal oscillation of the inter-layer phase difference due to the moving Josephson vortex lattice, with the frequency matching with that of any CTP modes, excites a resonant plasma oscillation. The Josephson vortex lattice also transforms its spatial configuration along the c axis in accordance with the c-axis standing-wave modes of the plasma oscillation. The resonance of Josephson vortex lattice to the CTP modes is revealed in the form of the multiple Josephson vortex-flow branches in the tunneling current-voltage (I-V) characteristics of a stack of IJJs.

Experimental observation of the multiple CTP modes in a stack of a finite number of IJJs has often been attempted in a mesa or an S-shaped stack of IJJs fabricated on the surface of a large stack of IJJs, which is called a pedestal or a basal layer. Since each superconducting bilayer in a stack of IJJs is much thinner than the c-axis London penetration depth λab, Josephson vortices in a usual mesa or an S-shaped-stack structure are in a strong inductive coupling to those in the basal layer that...
The Josephson vortex configuration in stacked IJJs of Bi-2212 single crystals is determined by the magnetic flux density in the junctions and the junction size. Josephson vortices are introduced to the insulating layer of a junction when its length \( L \) is longer than the Josephson penetration depth, \( \lambda_J = \sqrt{\Phi_0/2\pi\mu_0c}(D + 2\lambda_{ab}/s) \), where \( \Phi_0(=h/2e) \) is a flux quantum. One can identify various characteristic magnetic fields in relation with the dynamics of Josephson vortices [see Fig. 1]. For a stack of junctions of length \( L \) Josephson vortices start entering the junctions above a characteristic field \( H_{c1}^{||} = (\Phi_0\lambda_J/L^2\lambda_{ab}\ln(\lambda_{ab}/D)) \). This field can be higher than the nominal value of one flux quantum threading a junction of length \( L \), that is \( H_{c0}^{||} = (\Phi_0/(s + D)L) \), especially for a junction comparable to \( \lambda_J \). Even above \( H_{c0}^{||} \) the tunneling screening current prevents an external magnetic field from entering the junctions until it reaches \( H_{c1}^{||} \). Thus, for a junction with its length in the range of \( \lambda_J < L < L_m \), one can observe the usual oscillatory Fraunhofer modulation of the tunneling critical current as a function of an external magnetic field of \( H < H_{c1}^{||} \), although the junction size is larger than \( \lambda_J \) (Region 0).

The Josephson vortices entering the junctions for \( H > H_{c1}^{||} \) form a regular lattice as the magnetic field reaches the value \( H_{c2}^{||} = (1.4\Phi_0/2\pi(s + D)\lambda_J) \). In this field range (Region 1), since the field-induced vortices exist in arbitrary configurations in a given magnetic field, disordered multiple vortex-flow subbranches are generated in the \( I-V \) characteristics. In this region, the number of vortices in a junction and the propagation velocities of different vortex configurations are not correlated with each other. Thus, in this field range, the vortex-flow subbranches in the \( I-V \) characteristics are not related to vortex mode velocities or geometrical resonances.

In an external field above \( H_{c2}^{||} \) (Region 2), Josephson vortices start forming a triangular lattice. Recent observation of the oscillatory vortex-flow resistance arising from the interaction between the vortex lattice and the crystal boundary potential clearly confirmed the existence of the triangular vortex lattice above \( H_{c2}^{||} \). On the other hand, the resonance between the Josephson vortex lattice and CTP modes is well known in this vortex dynamics system. The collective resonance revealed as voltage peaks (in the case of the voltage bias) or voltage jumps (in the case of the current bias) has been searched extensively both theoretically and experimentally. The collective resonance has been observed in the dense-vortex state corresponding to the field range of \( H > H_{c3}^{||} = (\Phi_0/2\lambda_J(s + D)) \), where the shortest inter-vortex spacing becomes comparable to the diameter of a Josephson vortex 2\( \lambda_J \) (Region 3). Thus, Josephson vortices in Regions 2 and 3, although they form triangular lattices in both regions in a static state.
without a tunneling bias current (Fig. 1), behave differently in the vortex-flow state in the presence of a finite tunneling bias current. In the dynamic state, Josephson vortices in Region 2 remain in a triangular lattice, while those in Region 3 transform their lattice configuration to fit the plasma propagation modes in the stacked junctions.

The non-Josephson-like emission from the Josephson vortex-motion and the Shapiro resonance steps in the Josephson vortex-flow state have been observed in this dense-vortex state. These results confirm the existence of the coherent motion of the Josephson vortex lattice through the entire thickness of stacked (intrinsically Josephson junctions.

For the highest characteristic field \( H_{c1} \parallel \left[ \Phi_0 / \lambda_J (s + D) \right] \text{ (~4 T for Bi-2212) } \), the shortest inter-vortex spacing becomes comparable to the Josephson penetration depth, \( \lambda_J \). In this region (Region 4), the Josephson current along the length of a junction distributes almost sinusoidally so that the collective resonance behavior between the vortex motion and the plasma modes is expected to be stronger than in a field lower than \( H_{c1} \parallel \).

III. EXPERIMENT

As-grown slightly overdoped Bi-2212 single crystals were prepared by the conventional self-flux method. We fabricated, using the double-side cleaving technique, three samples of IJJs sandwiched between two Au electrodes at the top and the bottom of a stack of IJJs without the basal part [Fig. 2(b)]. This structure, which is in contrast to the usual mesa or the S-shaped-stack structure fabricated on the surface of a single crystal with a large basal layer, is more coincident with the model systems usually adopted in the theoretical analysis. The geometry with the basal layer eliminated by the double-side cleaving enables one to investigate the Josephson vortex dynamics in coupled IJJs without the interference from the vortex motion in the basal layer. Furthermore, the top and the bottom Au-film leads in a stack help the formation of the standing plasma oscillation modes along the c axis, while forming a transmission guide of the plasma oscillations along the length of the stack.

The samples containing sandwiched stacks were prepared in the following way. A single crystal was first glued on a sapphire substrate using negative photoresist and was cleaved until an optically clean and flat surface was obtained. Then, a 100-nm-thick Au film was thermally deposited on top of the crystal to protect the surface of the crystal and to make a good ohmic contact. A mesa structure was then prepared by photolithographic patterning and Ar-ion-beam etching. The surface of the patterned mesa was fixed to another sapphire substrate using negative photoresist and the basal layer of the mesa was subsequently cleaved away. This process, called the double-side-cleaving technique, was developed by Wang et al.. A 100-nm-thick Au film was again deposited on this freshly cleaved crystal surface, leaving a stack of IJJs sandwiched between two Au electrodes. A few-\( \mu m \)-long portions on both sides of the sandwiched stack were etched away subsequently [see Fig. 2(a)] to get the bottom Au electrode exposed for c-axis transport measurements as schematically depicted in Fig. 2(b). Finally, a 300-nm-thick Au-extension pad was attached on each of the Au electrode. The lateral sizes of three samples were \( 13.5 \times 1.4 \mu m^2 \) [SP1], \( 15 \times 1.4 \mu m^2 \) [SP2] and \( 16 \times 1.5 \mu m^2 \) [SP3].

The magnetic field was aligned in parallel with the plane of junctions within the resolution of 0.01 degree to minimize the pinning of Josephson vortices by the pancake vortices that can form in \( \text{CuO}_2 \) bilayers due to the vertical component of a misaligned field. The alignment was done, in a field of 4 T and at a temperature around 60 K, by tuning to the angle that gives the maximum Josephson-vortex-flow resistance or the minimum pinning of the Josephson vortices by the pancake vortices generated by the angle misalignment. To minimize the external noise, transport measurements were performed with a low-pass filter connected to each electrode located at room temperature.

![FIG. 2: (color online) (a) An optical micrograph of a sample during the fabrication process. The scale bar represents the length of 10 \( \mu m \). (b) Schematics of sample and measurement configurations.](image)

| sample number | stack size (\( \mu m^2 \)) | \( \lambda_J \) (\( \mu m \)) | \( H_{c1} \parallel \) (T) | \( c_m \) (\( \times 10^5 \) m/s) | \( c_1 \) (\( \times 10^5 \) m/s) |
|---------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| SP1           | 13.5×1.4        | 35              | 0.25            | 2.8             | 3.8             |
| SP2           | 15×1.4          | 22              | 0.32            | 2.1             | 3.9             |
| SP3           | 16×1.5          | 45              | 0.24            | 2.9             | 4.2             |

TABLE I: Sample parameters. \( N \) is the total number of intrinsic Josephson junctions in a stack, \( \lambda_J \) the Josephson penetration depth, \( H_{c1} \parallel = \Phi_0 / 2 \lambda_J (s + D) \) (s and D the thickness of the superconducting and the insulating layers, respectively), \( c_m \) the maximum propagation velocity of vortices in Region 3, and \( c_1 \) the velocity of the fastest collective plasma mode estimated by the RCSJ model.
IV. RESULTS AND DISCUSSION

A. Static and dynamic state

Fig. 3(a) shows the oscillation of Josephson-vortex-flow resistance of SP1 \(|N=35|\) as a function of external magnetic fields at \(T=55\) K in bias currents from 1 \(\mu A\) to 10 \(\mu A\) at intervals of 1 \(\mu A\). The field period of the oscillation, \(H_p=510\) G is one half of \(H_0 (=\Phi_0/L(s+D))\) through the junction area of length \(L=13.5\ \mu m\) and thickness \(s+D=1.5\ nm\). The oscillation of the Josephson-vortex-flow resistance is known to be caused by the interaction between the triangular Josephson vortex lattice along the c axis and the vortex-flow boundary potential at the junction edges.\(^{20,21,29}\) The observation of the Josephson-vortex-flow resistance oscillation in our sandwiched stack without the basal layers indicates that, as the S-shaped-stack geometry, our sandwiched geometry also supports the Josephson-vortex triangular lattice in the steady state [see the inset of Fig. 4(a)]. The oscillation in Fig. 3(a) is reduced by increasing bias current as the bias current tilts the boundary potential.\(^{33}\) The oscillatory-magnetoresistance behavior is also suppressed completely at \(T \sim 80\) K near \(T_c (=92\) K) as in the inset of Fig. 3(b), which is related to melting of the near-static Josephson vortex lattice.\(^{31}\)

The inset of Fig. 3(a) shows the single Josephson-vortex-flow branch (the low-bias region) and quasiparticle branches (the high-bias region) in \(H=1.2\) T at \(T=55\) K. The field, belonging to Region 2 in Fig. 1, corresponds to the position indicated by the arrow in the main panel. In this field region, the single-curve Josephson-vortex-flow branch caused by the triangular lattice persists even down to 4.2 K (not shown). In Fig. 3(b), I-V characteristics in a field of Region 3 at \(T=25\) K also show a single Josephson vortex-flow branch corresponding to the triangular lattice. Table I shows the characteristic field value of \(H_{13}^c\) of each sample.

Unlike the field range of Region 2, however, the single Josephson-vortex-flow branch in Region 3 evolves into the multiple Josephson-vortex-flow sub-branches with decreasing temperature. Fig. 3(b) shows the temperature dependence of the Josephson-vortex-flow branches, where the multiple sub-branches start appearing at \(T \sim 20\) K and become clearer with decreasing temperature. At \(T=4.2\) K, a single branch representing the triangular lattice below the current bias of \(I_b \sim 2\ \mu A\) splits into multiple Josephson-vortex-flow sub-branches for the higher current bias. The Josephson vortex flow, then, transits to the McCumber quasiparticle tunneling for \(I_b\) above \(8\ \mu A\). The multiple Josephson-vortex-flow sub-branches rapidly suppress with increasing temperature above 4.2 K and disappear completely as the temperature reaches 25 K. On the contrary, the quasiparticle-tunneling branches are robust to the temperature variation and persist up to \(T=25\) K. This contrasting temperature dependencies of two different regions strongly imply that the multiple branches in the two regions are of different physical origins.

We also observed the separation of the two characteristic multiple-branch regions in SP2 and SP3. Figs. 4(a) and 4(b) show the magnetic field dependencies of I-V characteristics of two sandwiched stacks, SP2 and SP3, respectively, at 4.2 K. The number of zero-field quasiparticle-tunneling branches of SP2 and SP3 were 22 and 45, respectively (not shown), which coincided with the total number of IJJ layers in the given stacks. The differentiation of the multiple branches by the collective Josephson vortex flow (the low-bias region) from the ones by the quasiparticle tunneling (the high-bias region) is based on the fact that the delimiting current, represented by the current values marked by the dotted vertical lines, stemmed from the zero-field Josephson critical current. Nonetheless, a finite resistance is present in the region.
dynamics of stacked Josephson junctions. Since the speed of vortices, \( c_m \), is limited by that of the maximum electromagnetic wave in a junction, Josephson-vortex-flow state changes to quasiparticle tunneling state near the maximum cut-off velocity or equivalently by the limiting voltage \( V_m \) in the outermost sub-branch. Using the relation \( c_m = \frac{H}{2}\pi \frac{I_c}{2} \) of \( V_m = NH(s + D)_j c_m \), we could get \( c_m \) as \( \sim 3.9 \times 10^5 \text{ m/s} \) and \( \sim 4.2 \times 10^5 \text{ m/s} \) for the corresponding delimiting bias point \( V_m \) in SP2 and SP3, respectively. These values are similar to those obtained from a Josephson-vortex-flow branch in mesa structures. In the case of the usual mesa or the S-shaped-stack structure with a basal layer, only a single Josephson-vortex-flow branch used to be observed below the critical bias point, regardless of the temperature and the magnetic field intensity. The velocity \( c_m \) corresponding to the delimiting value on this single branch has been regarded as the slowest mode velocity, \( n=N \). As illustrated below, however, in our case without the basal layer \( c_m \), corresponding to \( V_m \) in the outermost sub-branch, represents the propagation velocity of the \( n=1 \) plasma mode.

In the resistively and capacitively shunted junction (RCSJ) model, the junction capacitance is obtained by the relation, \( C_j = \beta c \frac{\Phi_0}{2\pi R_n} \), where \( \beta c \) \( = (4|I_c|/\pi I_s)^2 \) is the McCumber parameter, \( R_n \) is the normal-state junction resistance and \( I_s \) is the return current. In SP2, \( C_j \) is obtained to be about 47 pF, with \( I_c(=0.2 \text{ mA}), I_s(=4.8 \mu\text{A}) \) and \( R_n(=10 \Omega) \). The Swihart velocity \( c_0 = \sqrt{sA/2\mu_0C_jA_{nb}^2} \), estimated from the junction capacitance, is \( 3.6 \times 10^5 \text{ m/s} \), where \( A_{nb}=200 \text{ nm} \) and the junction area of \( A=21 \mu\text{m}^2 \). The resulting fastest mode velocity \( c_1 = 3.7 \times 10^5 \text{ m/s} \) is in remarkable agreement with the observed value of \( c_m \) in SP2. Experimental values of \( c_m \) in Table I are consistent with the calculated ones of \( c_1 \) for all three samples. This indicates that the outermost sub-branch is the fastest velocity, \( n=1 \), mode.

In general, Josephson-vortex resonance branches by the inter-layer inductive coupling have been predicted to generate current steps near the resonance voltages on the background of a single resistive branch. In the present study, however, Josephson-vortex resonance to the CTP modes constitutes multiple sub-branches over almost the entire bias range. A recent theory for an interlayer inductive coupling hybridized by a capacitive coupling predicts that the resistive sub-branches corresponding to the collective resonance modes show linear \( I-V \) characteristics in the low-bias region with the varied slopes from one another. More spread in the slopes are predicted for a higher capacitive coupling. Fig. 5(a) shows the Josephson-vortex-flow sub-branches for SP2 at \( H=5 \text{ T} \). As predicted by the theory, all the sub-branches in the Josephson vortex-flow regime exhibit linearities with varied slopes in the low bias range. The extrapolation of the linear low-bias region converges to a single point on the current axis, although the inductive-capacitive hybrid coupling theory predicts they should converge to the origin. The discrepancy indicates a finite pinning of Josephson vortex motion due to the pres-
ere the field. Here tunneling critical Josephson current at a given applied bias region. The best-fit value of \( \alpha \) illustrates the best fit to the sub-branches in the low-denoted with an arrow in Fig. 5(a). The gray region is the depinning current of the Josephson-vortex, \( I_c \) is the capacitive coupling represented by the parameter, \( \alpha \)

\[
\alpha = \left[ \frac{I_c - I_p}{V_n \cos(n\pi/N)} \right] + \Delta I_p - \Delta V_n
\]

where \( n=1,2,...,N \). Here, \( I_p \) is the depinning current of the Josephson-vortex, \( I_c \) is the tunneling critical Josephson current at a given applied field. Here \( V_n \) is the maximum mode voltage corresponding to a given bias current as represented by the curve denoted with an arrow in Fig. 5(a). The gray region illustrates the best fit to the sub-branches in the low-bias region. The best-fit value of \( \alpha \) turns out to be 0.45 with \( V_c = 4.22 \) mV. The best-fit values of \( \alpha \) were found to be almost insensitive to magnetic fields as \( \alpha = 0.43 - 0.45 \) for \( H = 4 - 5 \) T [see the inset of Fig. 5(a)], which is consistent with the assumption of the capacitive coupling and reasonably close to the theoretical expectation of \( 0.1 < \alpha < 0.4 \). Therefore, as shown in the upper inset of Fig. 4(a), we consider that the first \( (n=N) \); the leftmost and last \( (n=1) \); the rightmost) sub-branches correspond to (i) the triangular and (ii) the rectangular Josephson vortex lattice along the \( c \) axis, respectively.

Fig. 5(b) shows the change of Josephson-vortex-flow sub-branches for SP2 when magnetic field increases from \( H = 4.35 \) T to 4.5 T. The more changes take place in the voltages between the two magnetic fields for branches in the higher voltages (branches more to the right). This behavior can be compared to the dispersive relation of Eq. (1), which states that the lower-index modes are more dispersive than the higher-index modes. Thus, the right-most branches among multiple Josephson-vortex sub-branches in Figs. 3–6 are again considered to represent the square vortex lattice [see the inset of Fig. 4(a)] corresponding to the \( n = 1 \) plasma excitation mode.

The maximum number of vortex-flow sub-branches counted for various fields in SP2 is 20, which is two modes smaller than the number of quasiparticle branches (or, equivalently, than the number of junctions). Recently, Ryndyk proposed that the additional dissipation due to the charge-imbalance relaxation due to quasiparticle injection in a junction can prevent the structural transformation of Josephson-vortex lattice along the \( c \) axis. In our sample geometry quasiparticles are injected directly from outside through the outermost two junctions, which may cause the non-equilibrium charge-imbalance effect in the two outermost \( \text{CuO}_2 \) bilayer. The resulting suppression of the configurational transformation of Josephson-vortex lattice reduces the number of observable resonance modes. For inner \( \text{CuO}_2 \) planes, on the other hand, with the pair tunneling for a bias below the critical current of the junction, the quasiparticle injection can be neglected.

B. Temperature dependence

Figs. 6(a) and (b) show the temperature dependencies of \( I-V \) characteristics for SP2 in \( H = 5.2 \) T and for SP3 in \( H = 3 \) T, respectively. The temperature dependencies of the sub-branches in the Josephson vortex-flow region and quasiparticle-tunneling region are much different from each other. Although not clearly illustrated in the figures the quasiparticle branches survived up to temperatures close to \( T_c \). The sub-branches in the Josephson vortex-flow region, however, disappeared at a characteristic temperature \( T_{eh} \), near 18 K and 28 K for SP2 and SP3, respectively, leaving a single branch as SP1 in Fig. 3(b). These characteristic temperatures are almost insensitive to the fields applied [see insets of Figs. 6(a) and (b)].

In the vortex flow state for the magnetic-field range of Regions 3 and 4 the triangular vortex configuration and the multiple-mode configuration resonating with the col-
The Josephson vortex dynamics was comprehensively studied in compactly stacked intrinsic Josephson junctions over a wide range of magnetic fields and temperatures. Measurements were made on three stacks of intrinsic Josephson junctions, each sandwiched between two (top and bottom) normal-metallic electrodes, thus eliminating the interference from the basal layer(s). The $I$-$V$ characteristics for high external magnetic fields (Regions 3 and 4 defined in Fig. 1) showed three states: (i) the near-static triangular Josephson-vortex state with the formation of a single branch, (ii) the dynamic vortex state with multiple-mode sub-branches, and (iii) the McCumber quasiparticle-tunneling state with quasiparticle sub-branches. In the dynamic vortex state, each sub-branch corresponded to the different configuration of the Josephson-vortex lattice along the $c$ axis from a triangular to a rectangular lattice. With increasing temperature, the sub-branches disappeared at $T_{ch}$, above which no configurational transformation of the Josephson-vortex lattice can be assumed.

Due to the strong inductive coupling to many junctions in the basal layer(s), the mesa or the S-shaped-stack structure, although containing a finite number of junctions in the stack under study, behaves effectively as consisting of an infinite number of junctions along the $c$ axis. Thus, these structures cannot support the formation of the multiple Josephson-vortex-flow modes as predicted theoretically for a finite number of junctions. Although the $c$-axis boundary condition is not strictly imposed in these systems the transverse vortex motion is still restricted by the boundary conditions at the side edges. This is the main reason why, in the mesa or the S-shaped-stack structure, clear magnetoresistance oscillations are observed by the interplay between the near-static triangular Josephson-vortex configuration and the boundary potential in the systems, while the transverse Josephson-vortex multiple modes are not obtained. One should be noted that the multiple collective Josephson vortex modes can be obtained only in a stack containing a finite number of (intrinsic) Josephson junctions. The triumph of this study is that our samples truly satisfied this condition of a finite number of junctions in a stack and, as a result, both the near-static Josephson vortex state highly coupled to the edge potential and the collective transverse Josephson-vortex state were identified in each junction. This facilitated the detailed identification of the vortex configurations illustrated in Fig. 1.

V. CONCLUSIONS

FIG. 6: Temperature dependence of $I$-$V$ characteristics for (a) SP2 in $H=5.2$ T and (b) for SP3 in $H=3$ T. The curves are shifted vertically for clarity. The insets of (a) and (b) show a single Josephson vortex-flow branch above the characteristic temperature $T_{ch}$ in two different magnetic fields for SP2 and SP3, respectively.
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