Anomalous Josephson Effect in Andreev Interferometers

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The first evidence of anomalous Josephson effect in mesoscopic superconductor-normal metal-superconductor (SNS) junctions forming a crosslike Andreev Interferometer is reported. From conductance measurements of the out-of-equilibrium weak link, we determine a voltage-controlled spontaneous phase that resembles a \(\varphi_0\)-junction in the absence of magnetic and spin-orbit interactions normally required to break time-reversal and inversion symmetries. The temperature and voltage dependences together with a dissipative term in the current-phase relation strongly suggest that the mechanism underlying the observed phenomenon relies on electron-hole asymmetries.

The dc Josephson effect establishes that current flows without dissipation across two superconductors interrupted by a weak link in a so-called Josephson junction (JJ) [1]. The supercurrent \(I_s\) and the macroscopic phase difference of the superconductors \(\delta\) are correlated by the current-phase relation (CPR) which hosts some general properties [2, 3] like \(2\pi\) periodicity and time-reversal symmetry \(I_s(-\delta) = -I_s(\delta)\). The latter determines the \(I_s(\delta = 0) = 0\) condition and a CPR given by a series of sinusoidal functions. However, most of these junctions are well described by the first harmonic \(I_s(\delta) = I_s\sin\delta\) with minimum energy at \(\delta = 0\) and therefore called 0-JJs. Under certain conditions, the supercurrent can change its polarity leading to a \(\pi\) ground state. Discovered in nonequilibrium-controlled metallic JJs [4, 5], \(\pi\)-JJs were intensely investigated thereafter in ferromagnetic [6, 7], semiconductor-based JJs [8, 9] or using superconductors with unconventional paring symmetry [10, 11]. However, these JJs still present a rigid ground state given by the supercurrent direction which constrains the flowing Josephson current to \(\delta \neq n\pi\) states with \(n\) any integer. To date, two mechanisms have been explored to achieve JJs with variable ground states. Those with a large and negative second harmonic that gives rise to a degenerate and arbitrary design-determined \(\pm\varphi\) phase [12–14]; and junctions with broken inversion symmetry and a non-degenerate and controllable \(\varphi_0\) phase bias acting as a Josephson phase battery [15–18]. The former is possible in a combination of \(0-\pi\)-JJs owing to spatial oscillations of the Cooper pair wave function [19, 20], ferromagnetic weak links are the most suitable candidates thanks to the ground state dependence on the magnetic layer thickness [21, 22]. In \(\varphi_0\)-JJs, the interplay of an exchange splitting field and strong spin-orbit coupling (SOC) [23–25] or noncoplanar magnetic texture [26] induces a finite phase shift

\[
I_s = I_c \sin(\delta - \varphi_0) = I_J \sin\delta + I_{an} \cos\delta
\]

with \(I_J = I_c \cos \varphi_0\) and \(I_{an} = -I_c \sin \varphi_0\) the usual and anomalous Josephson current. In diffusive metallic junctions, the amplitude and direction of the supercurrent is given by the occupation of the supercurrent-carrying density of states (SCS) travelling parallel to the phase gradient and those carrying the supercurrent in the opposite direction [27, 28]. Unlike previous systems, \(0-\pi\) transition is achieved by driving the electrons out-of-equilibrium. A mesoscopic conductor coupled to the weak link in the so-called Andreev Interferometers provides the nonthermal distribution [29] that allows the selective depopulation of the dominant positive states. A recent revision of the quasiclassical Green’s function formalism suggested that electron-hole asymmetries may induce non-vanishing Aharonov-Bohm-like currents in multiterminal normal-superconducting heterostructures [30–32]. The interplay of even-in-\(\delta\) Aharonov-Bohm oscillations and normal Josephson currents could be responsible for \(\varphi_0\)-states in asymmetric interferometers. However, no such states have been reported so far.

In this letter, we report on the demonstration of metallic and non-magnetic \(\varphi_0\)-junctions forming mesoscopic crosslike Andreev interferometers known as Hybrid Quantum Interference Device (HyQUID) [33]. By embedding the anomalous junction in a superconducting ring with non-negligible screening currents, the resistance of the proximitized normal wire connected to the JJs enters in a nonlinear regime with bistable states and a highly \(\varphi_0\)-dependent hysteretic behavior. The supercurrent in our lateral hybrid junctions is suppressed by modifying the electrostatic potential of the weak link \(V_N\) and voltage-controlled \(\varphi_0\)-states emerge beyond a critical value.

We fabricated Nb/Ag/Nb HyQUIDs depicted in Fig. 1a,b which consist of two 300 nm-thick superconducting electrodes interrupted by the non-magnetic 50 nm-thick metal connected to reservoirs in a crosslike geometry (dark yellow). Our junctions exhibit irregular but highly transparent S/N interfaces with S-electrodes spaced \(L_x = 500\) nm larger than the superconducting \(\xi_0\) and the normal \(\xi_N\) coherence length which places our devices in the diffusive and long junction regime. By applying a voltage across the connecting normal wires with to-
FIG. 1. (a) HyQUID schematic with the 4-wire electrical configuration for measuring the interferometer (dark yellow) differential resistance $R_N$. (b) False-colored micrograph of the junction (superconducting electrodes in blue). (c) Flux-tilted washboard potential of the system for $\phi_0 = -\pi/4$ and $\beta = 10$. Positive (negative) external flux shifts to the left (right) the $\Phi_0$-periodic potential and the particle escapes to a lower energy state at $\Phi_0 + n(\Phi_0 - n)$, red (green) balls. (d) $\delta(\Phi_e)$ for $\phi_0 = 0, -\pi/4, \pi/4$. with phase jumps at $d\Phi_e/d\delta = 0$. (e) Colormap plot of a HyQUID differential resistance for positively sweep flux and bias voltages $V < V_{c1}$. (e) $R_N(\Phi_e)$ traces for selected voltages with positive (red) and negative (green) flux polarity. $R_N^+(\Phi_e^+) > R_N^-(\Phi_e^-)$ for all the bias voltages. The gray curves are the best fit to the model. Phase jump transitions are included in the dashed lines.

tallength $L_N \simeq 4$ $\mu$m, we created a nonthermal quasiparticle energy distribution function. Large reservoirs and a diffusion time $\tau_D = L_N^2/D$ shorter than the characteristic interaction time ensure out-of-equilibrium conditions. A diffusion coefficient $D = 150$ cm$^2$/s and a Thouless energy $E_{th} = \hbar D/L_N^2 = 0.04$ meV are estimated from resistance measurements (see details in Ref. [34]). We conduct measurements on two different configurations of HyQUID, one in which the superconducting electrodes are connected to a four-wire electrical setup in a similar fashion than in Ref. [4, 5], allowing the direct probing of the critical current, and one in which the electrodes are joined to form a superconducting loop with an inductance $L = 290$ pH. The latter allows to control the phase across the junction by the application of a magnetic flux and determine the anomalous phase independently of any spurious flux trapped in the system [35]. If the screening parameter $\beta = 2\pi L I_c/\Phi_0 > 1$, the contribution to the net flux due to the supercurrent-induced flux results in a phase difference that is a non-linear function of the external flux $\Phi_e$:

$$\delta = 2\pi \frac{\Phi_e}{\Phi_0} - \beta i = \phi_c - \beta i$$  \hspace{1cm} (2)

with $i = I_s/I_c$ the normalized supercurrent and $\phi_c \equiv 2\pi \Phi_e/\Phi_0$ the “external phase”. The potential energy for such junctions is given by the Josephson energy and the magnetic energy stored in the loop inductance [36]

$$U = E_J \left[ 1 - \cos(\delta - \varphi_0) + \frac{(\delta - \phi_c)^2}{2\beta} \right]$$  \hspace{1cm} (3)

with $E_J = \Phi_0 I_c/2\pi$. If $\beta < 1$, $\delta \simeq \phi_c$ and the potential energy has a parabolic shape with a single minimum. For $\beta > 1$, the oscillating component gives rise to several minima and maxima according to the positive and negative slopes of the multi-valued function $\delta(\phi_c)$. $U(\delta)$ presents inflection points at

$$\phi_{n\pm} = \varphi_0 \pm [\phi_t(\beta) + 2\pi n]$$  \hspace{1cm} (4)

at which $\delta(\phi_c)$ slopes change sign. $\phi_t = \sin^{-1}[1/\beta] + \sqrt{\beta^2 - 1 + \pi/2}$, see calculations in Ref. [34], varies linearly for large $\beta$. For convenience, we use $\Phi \equiv \phi_c/2\pi$ notation. Under the classical analogy of a particle moving in a “washboard potential”, the particle is trapped in a potential minimum $U = 0$. When an external flux is applied, the parabolic part is shifted along the oscillating component and the particle energy increases captured in the potential well. The well barriers decrease with the flux
FIG. 2. (a) Colormap plot of the HyQUID differential resistance for positive (left) and negative (right) sweeping flux polarity above $V_{c1}^N$. The evolution is described by zones (dashed squares) in text. (b) Characteristic $R^+_{N}$ (red) and $R^-_{N}$ (green) traces for each zone at $V_N = 0.33, 0.4, 0.56$ and $0.67$ mV. (c) Effective critical current $|I_c|$ and $\varphi_0$-states as a function of the bias voltage. The usual Josephson current $I_J$ reverses at $V^*_N \approx 0.44$ mV with a jump discontinuity according to the model. The system eventually tends to a state dominated by the anomalous Josephson current with $|I_c| \to 0$ and $\varphi_0 = -\pi/2$. until vanishing at $\Phi_n = \Phi_{n=0}^\pm$, the induced supercurrent $I_s = \Phi_n/L$ matches the critical current, a flux quantum enters the loop reducing $|\phi_e - \delta|$, and the particle escapes to a lower energy state in a “phase jump”. In $\varphi_0$-JJs, the zero-flux minimum shifts to $\delta = \varphi_0$. The potential remains invariant with respect to the simultaneous change of $\phi_e$ and $\delta$ by $2\pi n$, but the anomalous CPR breaks the phase difference symmetry $\delta(\Phi_e) \neq -\delta(-\Phi_e)$ and the system becomes direction-dependent $\phi^+ = \phi^- + 2\varphi_0$ (Fig. 1c,d).

The differential resistance $R_N$ of the coupled connector was measured at $T = 2.7$ K as a function of the inductively applied flux swept back and forth and the voltage at the center of the wire given by $V_N = I_N R_N(0)/2$. $I_N$ is the bias current and $R_N(0)$ the differential resistance at zero bias current. All devices exhibit similar resistance $R_N(0) = 1.8(2)\Omega$. The conductance oscillates as a function of the phase difference due to phase transfer from the superconducting condensate to normal electrons via Andreev reflections at the $S/N$ interfaces [37] with a minimum resistance value at $\delta = 2n\pi$ and equals the normal-state resistance $R_N^0$ at $\delta = (2n+1)\pi$ [38], thereby it can be approximated by

$$R_N = R^0_N - \frac{\Delta R}{2}[1 + \cos \delta] \quad (5)$$

a quantitative study of $\Delta R(V_N)$ is given in Ref. [34]. $R_N(\Phi_e, V_N)$ colormap and selected traces are shown in Fig. 1ef for $V_N < V_{c1}^N = 0.26$ mV. Gray curves are the best fit to the model given by Eq. 1, 2 and 5 with the phase jumps (dashed trajectories) obtained from $\Phi_{n=0}^\pm$ and $\Delta R, \beta, \varphi_0$ the fitting parameters. $\Phi_{n=0}$ decreases monotonically until the critical current is suppressed at $eV_{c1}^N \simeq 7E_{th}$ where the system becomes fully periodic. Resistance at phase jumps is slightly lower when the flux is negatively swept $R^+_N(\Phi_e^+) > R^-_N(\Phi_e^-)$, explained by a small and quasi-constant anomalous current $I_{an} < 0$ corresponding to $\varphi_0 = +\pi/4(2 \pm 1)$. For large $\beta$, the resulting state of the phase jumps are not the one with the lowest energy. If barriers are overcome, e.g., by thermal fluctuations, several flux quanta may enter the loop and the system falls to a much lower energy or materialize at $\Phi^* < \Phi_n$. These phenomena are observed for almost any voltage regardless of the flux polarity. Notwithstanding the above, the model adequately tracks these anomalies as the system remains trapped in the next available potential well as the gray lines shown.
The reflection symmetry is fully broken above $V_N^{ \text{th}}$, (Fig. 2a,b). $R_N^c(\Phi_n)$ becomes non-periodic owing to screening currents $\beta>1$, but supercurrent is not reversed as expected in symmetric interferometers. $R_N$ shows a sawtooth pattern with jumps from high to low values for both flux polarities. As the bias voltage increases, the resistance jumps in $R_N^c$ do not drop to the lowest resistance state and the curves become hysteretic. The height of the resistance jumps progressively decreases in $R_N^c$ and the gap between the curves increases, well described by a monotonic increase of both the screening currents and a negative phase shift $0<|\varphi_0|<\pi/2$. The sign change in $\varphi_0$ means that only $I_{an}$ is reversed. At $V_N=0.38$ mV, $R_N^c$ jumps disappear at $\varphi_0=-\pi/4$ and $\beta \simeq 11$. Beyond this point (purple zone), $R_N^c$ reaches the minimum value and remains constant with some traces of small jumps from low to high resistance states, meanwhile $R_N$ jumps are progressively rounded. The model predicts small phase jumps in $R_N^c$ for larger $|\varphi_0|$ with a flip in the direction of the resistance jumps at $\varphi_0 \simeq -\pi/3$ and large screening currents. Above $V_N^{\text{th}} \approx 0.44$ mV, $R_N^c$ peaks disappear, the resistance increases smoothly and remains constant, signature of the $I_z$ inversion, $\pi>|\varphi_0|>\pi/2$. Phase jumps in $R_N^c$ are expected to disappear for $\varphi_0 = -l \pi/6$ and $\beta>5$ with $l \in 4-5$ as it is shown for $V_N=0.56$ mV. The onset of $R_N^c$ resistance jumps in the green zone with rounded maxima is explained by the decrease of $|\varphi_0|$ and the screening currents that pull the $R_N$ branches apart for $\beta<5$ and $\varphi_0 \to -2\pi/3$. $R_N^c$ recovers the sawtooth shape with a similar profile than $R_N^c$ for $V_N>0.63$ mV, consistent with a CPR controlled by $I_{an}$, $\varphi_0 \simeq -\pi/2$. The gap decreases and the system becomes periodic and fully symmetric $R_N^c=R_N^c$ for $V_N>0.68$ mV. For $|\varphi_0| = \pi$, $R_N^c$ is not tilted and the resistance branches cross near their maximum value at a given flux $\Phi^c(\beta)$. For $V_N \in V_N^{\text{th}}-0.5$ mV, we assumed that the escape rate from the potential well is non-negligible at $\Phi^c<\Phi_n$ and the model can fit the data considering the initial curvature and the $R_N^c$ jumps. The fitting parameters are plotted in Fig. 2c with a jump-like transition to quasi-$\pi$ states at $V_N^{\text{th}}$. Thermal activated phase jumps (TAPJ) around $V_N^{\text{th}}$ are discussed in Ref. [34].

We measured the differential resistance of the junction as a function of the bias voltage applied as before in current-biased interferometers. The I-V characteristics, obtained by integration, for selected $V_N$ values are shown in Figure 3a. $I_c$ decreases monotonically mimicking the results obtained in flux-biased HyQUIDs with a fully resistive state at a critical value $V_N^{\text{th}}$. The persistent current emerges for higher voltages until it vanishes definitively at $eV_N^{\text{th}} \approx 0.9$ mV $\approx 22 E_\text{th}$. A dissipative and voltage-dependent current $I_0$ displaces the superconducting region $I_s$ (gray area). $I_0$ scales linearly with the applied voltage up to $V_N \approx 0.7$ mV, beyond which thermal effects undermine phase-coherence correlations as shown in long junctions [34]. Measurements in sample B with reversed polarity $I_0(\mp V_N) < 0$ rules out any leakage currents from the control wires to the superconducting electrodes [39].

The temperature dependence of the critical voltage is compared with the existing theory of fully symmetric interferometers [27, 28]. The electrostatic potential de-populates states with $\epsilon<eV_N$. Temperature undermines the selectivity of the de-population of positive states due to the temperature-rounded staircase-like occupation probability [29], increasing $V_N^{\text{th}}$ until the supercurrent is no longer reversed. In our HyQUIDs, the critical voltage remains constant or decrease with temperature up to $T \geq 6$ K. The expected values are plotted together with the experimental values in Fig. 3c. Details of the calculations in Ref. [34]. The fact that the theoretical $V_N^{\text{th}} \approx V_N^{\text{th}}$ suggests that the anomalous current modifies the SCS and reduces the critical voltage.

Spin-orbit interactions and/or a non-coplanar spin structure can be ruled out as the driving force since we used non-magnetic metals and the characteristic Kondo upturn caused by magnetic impurities in $R_N(0,T)$ was not observed. Unconventional pairing symmetries are not expected in Nb-based devices and the evolution of the differential resistance cannot be fitted if higher harmonics are considered. We also considered the dimensionality of our systems. The devices presented here are two-dimensional structures since $\xi_N = \sqrt{\hbar D/2 \pi k_B T} \approx 83$ nm $< L_x = L_y$. We e-beam patterned the metal crosses in finer HyQUIDs with similar S/N interfaces and trans-
parency but narrower weak links and control wires $L_y \approx 100$ nm. AJE is still noticeable in these devices despite the smaller critical currents, details in Ref. [34]. Nevertheless, the dissipative $I_0(\psi)$ and Aharonov-Bohm-like $I_{AB}$ terms in the CPR were predicted in topologically-modified interferometers [31, 32] which might stem from irregular interfaces [40]. For $\epsilon V_N > E_{ih}$ and highly transparent interfaces $I_{AB}$ are expected to be comparable with $I_1$ resulting in $\varphi(N)$-junction-states. The exponential decay of the Josephson currents [41] with temperature weaker than the $I_{AB} \propto 1/T$ [42, 43] would explain the anomalous temperature dependence and $\pi/2$-states at large voltages [30].

In summary, our work reveals the first evidence of anomalous Josephson effect in metallic and non-magnetic SNS junctions configured as Andreev interferometers. Our model provides a route to unambiguously determine the current-phase relation if large screening currents are present. $\varphi_N$-states can be tuned by modifying the metal electrostatic potential and probed in flux-biased junctions. HyQUIDs do not require side gates to control the electron density in semiconductor-based $\varphi_N$-junctions, complex 2DEG systems or scarce unconventional superconductors.

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[1] B. Josephson, Phys. Lett. 1, 251 (1962).
[2] K. K. Likharev, Rev. Mod. Phys. 51, 101 (1979).
[3] A. A. Golubov, M. Y. Kupriyanov, and E. Il’ichev, Rev. Mod. Phys. 76, 411 (2004).
[4] J. A. Baselmans, A. F. Morpurgo, B. J. van Wees, and T. M. Klapwijk, Nature 397, 43 (1999).
[5] R. Shaiikhaïdarov, A. F. Volkov, H. Takayanagi, V. T. Petrashev, and P. Delsing, Phys. Rev. B 62, R14649 (2000).
[6] T. T. Heikkilä, F. K. Wilhelm, and G. Schön, EPL, 51, 434 (2000).
[7] V. V. Ryazanov, V. A. Oboznov, A. Y. Rusanov, A. V. Veretennikov, A. A. Golubov, and J. Aarts, Phys. Rev. Lett. 86, 2427 (2001).
[8] C. T. Ke, C. M. Moehle, F. K. de Vries, C. Thomas, S. Metti, C. R. Guinn, R. Kallaher, M. Lodari, G. Scappucci, T. Wang, R. E. Diaz, G. C. Gardner, M. J. Manfra, and S. Goswami, Nat. Commun. 10, 3764 (2019).
[9] J. A. van Dam, Y. V. Nazarov, E. P. A. M. Bakkers, S. De Franceschi, and L. P. Kouwenhoven, Nature 442, 667 (2006).
[10] R. R. Schulz, B. Chesca, B. Goetz, C. W. Schneider, A. Schmelz, H. Bielefedt, H. Hilgenkamp, J. Mannhart, and C. C. Tsuei, Appl. Phys. Lett. 76, 912 (2000).
[11] H. J. H. Smilde, Ariando, D. H. A. Blank, G. J. Gerritsma, H. Hilgenkamp, and H. Rogalla, Phys. Rev. Lett. 88, 057004 (2002).
[12] E. Goldobin, D. Koelle, R. Kleiner, and A. Buzdin, Phys. Rev. B 76, 224523 (2007).
[13] H. Sickinger, A. Lipman, M. Weides, R. G. Mints, H. Kohlstedt, D. Koelle, R. Kleiner, and E. Goldobin, Phys. Rev. Lett. 109, 107002 (2012).
[14] E. Goldobin, H. Sickinger, M. Weides, N. Ruppelt, H. Kohlstedt, R. Kleiner, and D. Koelle, Appl. Phys. Lett. 102, 242602 (2013).
[15] D. B. Szombati, S. Nadj-Perge, D. Car, S. R. Plissard, E. P. A. M. Bakkers, and L. P. Kouwenhoven, Nat. Phys. 12, 568 (2016).
[16] A. Assouline, C. Feuillet-Palma, N. Bergeal, T. Zhang, A. MottaghiZadeh, A. Zimmers, E. Lhuillier, M. Eddiere, P. Atkinson, M. Aprili, and H. Aubin, Nat. Commun. 10, 126 (2019).
[17] E. Strambini, A. Iorio, O. Durante, R. Citro, C. Sanz-Fernández, C. Guarcello, I. V. Tokatly, A. Braggio, M. Rocci, N. Ligato, V. Zannier, L. Sorba, F. S. Bergeret, and F. Giazotto, Nat. Nanotechnol. 15, 656 (2020).
[18] W. Mayer, M. C. DartiaIli, J. Yuan, K. S. Wickramasinghe, E. Rossi, and J. Shabani, Nat. Commun. 11, 212 (2020).
[19] R. G. Mints, Phys. Rev. B 57, R3221 (1998).
[20] A. Buzdin and A. E. Koshelev, Phys. Rev. B 67, 220504(R) (2003).
[21] M. Weides, M. Kemmler, H. Kohlstedt, R. Wasen, D. Koelle, R. Kleiner, and E. Goldobin, Phys. Rev. Lett. 97, 247001 (2006).
[22] C. Gürlich, S. Scharinger, M. Weides, H. Kohlstedt, R. G. Mints, E. Goldobin, D. Koelle, and R. Kleiner, Phys. Rev. B 81, 094502 (2010).
[23] A. Buzdin, Phys. Rev. Lett. 101, 107005 (2008).
[24] F. S. Bergeret and I. V. Tokatly, EPL 110, 57005 (2015).
[25] F. Dolcini, M. Houzet, and J. S. Meyer, Phys. Rev. B 92, 035428 (2015).
[26] M. A. Silaev, I. V. Tokatly, and F. S. Bergeret, Phys. Rev. B 95, 184508 (2017).
[27] F. K. Wilhelm, G. Schön, and A. D. Zaikin, Phys. Rev. Lett. 81, 1682 (1998).
[28] T. T. Heikkilä, J. Särkkä, and F. K. Wilhelm, Phys. Rev. B 66, 184513 (2002).
[29] H. Pothier, S. Guéron, N. O. Birge, D. Esteve, and M. H. Devoret, Phys. Rev. Lett. 79, 3490 (1997).
[30] P. E. Dolgirev, M. S. Kalenkov, and A. D. Zaikin, Phys. Rev. B 97, 054521 (2018).
[31] P. E. Dolgirev, M. S. Kalenkov, A. E. Tarkhov, and A. D. Zaikin, Phys. Rev. B 100, 054511 (2019).
[32] P. E. Dolgirev, M. S. Kalenkov, and A. D. Zaikin, Sci. Rep. 9, 1301 (2019).
[33] V. T. Petrashev and C. Checkley, US. Patent No. 10859641B2 (2020).
[34] Supplemental Information.
[35] C. Guarcello, R. Citro, O. Durante, F. S. Bergeret, A. Iorio, C. Sanz-Fernández, E. Strambini, F. Giazotto, and A. Braggio, Phys. Rev. Res. 2, 023165 (2020).
[36] K. K. Likharev, Gordon and Breach Publishers, Amsterdam (1986).
[37] V. T. Petrashev, V. N. Antonov, P. Delsing, and T. C. Claeson, Phys. Rev. Lett. 74, 5268 (1995).
[38] V. T. Petrashev, R. S. Shaiikhaïdarov, I. A. Sosnin, P. Delsing, T. C. Claeson, and A. Volkov, Phys. Rev. B 58, 15088 (1998).
[39] By changing $V_N$ polarity any stray current should flow to the junction ground since the voltage terminals are connected to high input impedance amplifiers so $I_0(−V_N)$ would still be positive ( ).
[40] Electron-hole symmetry can be broken by introducing asymmetries in Andreev interferometers. Although existing work considers geometrical asymmetries, no work has investigated S/N interfaces with different transparencies to the authors’ knowledge. HyQUIDs present irregular and uneven interfaces that could provide the required asymmetry (4).

[41] P. Dubos, H. Courtois, B. Pannetier, F. K. Wilhelm, A. D. Zaikin, and G. Schön, Phys. Rev. B 63, 064502 (2001).
[42] H. Courtois, P. Gandit, D. Mailly, and B. Pannetier, Phys. Rev. Lett. 76, 130 (1996).
[43] A. A. Golubov, F. K. Wilhelm, and A. D. Zaikin, Phys. Rev. B 55, 1123 (1997).