A macroscopic fluid pump works according to the law of Newtonian mechanics and transfers a large number of molecules per cycle (of the order of $10^{23}$). By contrast, a nano-scale charge pump can be thought as the ultimate miniaturization of a pump, with its operation being subject to quantum mechanics and with only few electrons or even fractions of electrons transferred per cycle. It generates a direct current in the absence of an applied voltage exploiting the time-dependence of some properties of a nano-scale conductor. The idea of pumping in nanostructures was discussed theoretically a few decades ago [1–4]. So far, nano-scale pumps have been realised only in systems exhibiting strong Coulombic effects [5–12], whereas evidence for pumping in the absence of Coulomb-blockade has been elusive. A pioneering experiment by Switkes et al. [13] evidenced the difficulty of modulating in time the properties of an open mesoscopic conductor at cryogenic temperatures without generating undesired bias voltages due to stray capacitances [14, 15]. One possible solution to this problem is to use the ac Josephson effect to induce periodically time-dependent Andreev-reflection amplitudes in a hybrid normal-superconducting system [16]. Here we report the experimental detection of charge flow in an unbiased InAs nanowire (NW) embedded in a superconducting quantum interference device (SQUID). In this system, pumping may occur via the cyclic modulation of the phase of the order parameter of different superconducting electrodes. The symmetry of the current with respect to the enclosed magnetic flux [17, 18] and bias SQUID current is a discriminating signature of pumping. Currents exceeding 20 pA are measured at 250 mK, and exhibit symmetries compatible with a pumping mechanism in this setup which realizes a Josephson quantum electron pump (JQEP).

The microscopic mechanism that enables the transport properties of the NW to be affected by the phases of the superconducting order parameter is Andreev reflection [19]. This is the quantum process for which an electron impinging from the normal side onto the interface between a normal metal and a superconductor, is retroreflected as a hole (i.e., a time-reversed electron) which picks up the phase of the superconducting order parameter. When two or more superconductors are connected to the NW, multiple Andreev scattering processes can occur between them so that transport through the NW will depend on the differences between the phases of the order parameters [20].
The physical realization of this scheme is shown in Fig. 1a and consists of a heavily-doped InAs semiconducting NW on top of which three fingers of superconducting (S) vanadium (V) are deposited thus implementing a SQUID [21]. Two Au normal-metal electrodes (N) are coupled to the ends of the NW to allow detection of the current \( I_{\text{wire}} \) flowing through the wire. A close-up of the device core is shown in Fig. 1b. Time-dependence, and possibly pumping, arises from biasing the loop with a current \( I_{\text{SQUID}} \) larger than the critical current \( I_c \) of the SQUID so that the phase differences \( \varphi_1(t) \) and \( \varphi_2(t) \) across the two Josephson junctions cycle in time at the Josephson frequency \( \nu_J = V_{\text{SQUID}}/\Phi_0 \), where \( V_{\text{SQUID}} \) is the voltage developed across the SQUID and \( \Phi_0 \approx 2 \times 10^{-15} \text{ Wb} \) is the flux quantum. In addition, \( \varphi_1(t) \) and \( \varphi_2(t) \) can be shifted by a constant term \( \delta \varphi = 2\pi \Phi/\Phi_0 \) originating from an applied magnetic flux \( \Phi \) threading the loop. This scheme has the advantage that no high-frequency signal needs to be brought to the sample thus simplifying the setup and minimizing the impact of stray capacitances: the time-dependent signal is self-generated thanks to the ac Josephson effect.

Below the critical temperature of the superconductors (\( T_c \approx 4.65 \text{ K} \)) a Josephson current flows through the SQUID across the NW. The SQUID voltage-current characteristics at 250 mK is shown in the inset of Fig. 1c for two representative values of \( \Phi \). Whereas for \( \Phi = 0 \) the characteristics at 250 mK is shown in the inset of Fig. 1c for the two Josephson junctions cycle in time at the Josephson frequency \( \nu_J = V_{\text{SQUID}}/\Phi_0 \), where \( V_{\text{SQUID}} \) is the voltage developed across the SQUID and \( \Phi_0 \approx 2 \times 10^{-15} \text{ Wb} \) is the flux quantum. In addition, \( \varphi_1(t) \) and \( \varphi_2(t) \) can be shifted by a constant term \( \delta \varphi = 2\pi \Phi/\Phi_0 \) originating from an applied magnetic flux \( \Phi \) threading the ring. When \( \Phi_0 \) exceeds the SQUID critical supercurrent the ac Josephson effect sets up inducing a current \( I_{\text{wire}} \) which flows in the NW. \( I_{\text{wire}} \) is sensed through an ammeter. S and N denote superconductors and normal metals, respectively. (b) Color plot of the SQUID flux-to-voltage transfer function \( V_{\text{SQUID}} = \partial V_{\text{SQUID}}/\partial \Phi \) versus \( \Phi \) and \( I_{\text{SQUID}} \). \( V_{\text{SQUID}} \) is antisymmetric in \( \Phi \) and \( I_{\text{SQUID}} \). (c) Color plot of the NW flux-to-current transfer function \( I_{\text{wire}} = \partial V_{\text{wire}}/\partial \Phi \) versus \( \Phi \) and \( I_{\text{SQUID}} \). Data are taken with a voltmeter in an open-circuit configuration, i.e., without allowing \( I_{\text{wire}} \) to flow. Note the markedly different behavior displayed by \( I_{\text{wire}} \) and \( V_{\text{wire}} \) which are almost symmetric in \( \Phi \) as well as in \( I_{\text{SQUID}} \). All measurements are taken at \( T = 250 \text{ mK} \) with low-frequency phase-sensitive technique to get higher sensitivity and reduced noise.

\[
V_{\text{wire}} = \partial V_{\text{wire}}/\partial \Phi \quad \text{(Fig. 2d), where} \quad V_{\text{wire}} \quad \text{is measured with open NW contacts.} \quad I_{\text{wire}} \quad \text{and} \quad V_{\text{wire}} \quad \text{result from different but complementary measurements, and the evidence of such a similarity suggests that both reflect the same physical mechanism (see Supplementary Information).} \quad \text{As we shall argue, the nature of the symmetries displayed by} \quad I_{\text{wire}} \quad \text{and} \quad V_{\text{wire}} \quad \text{is compatible with a quantum pumping mechanisms.}}
\]

In general, the pumped current is not expected to show definite parity with \( \Phi \) [17, 18], therefore \( I_{\text{wire}} \) can have a flux-symmetric component as well. This, however, could be ascribed also to other mechanisms than pumping. In
addition, $I_{wire}$ is even not expected to possess any definite parity with $I_{SQUID}$. In order to extract a pure pumped current contribution from the whole measured signal we focus on the component of $I_{wire}$ which is antisymmetric in $\Phi$, $I_{wire}^{A}$, as it is predicted to be a fingerprint of quantum pumping in the JQEP [16]. After $\Phi$-integration of $I_{wire}$, $I_{wire}^{A}$ is therefore obtained as $I_{wire}^{A} = [I_{wire}(\Phi, I_{SQUID}) - I_{wire}(-\Phi, I_{SQUID})]/2$. The result of this procedure is shown in Fig. 3a which displays $I_{wire}^{A}$ versus $\Phi$ and $I_{SQUID}$ at 250 mK. The $\Phi_{0}$ periodicity joined with the antisymmetry imply that $I_{wire}^{A}$ vanishes at $\Phi = \Phi_{0}/2$, while its sign and magnitude can be changed by varying $\Phi$. Notably, $I_{wire}^{A}$ is almost symmetric in $I_{SQUID}$. The theoretical $I_{wire}^{A}$ calculated for the JQEP geometry through a dynamical scattering approach [23] assuming for the NW multiple independent modes is shown in Fig. 3b (see Supplementary Information). Although rather idealized, the model is an essential tool to predict the pumped current symmetries of the JQEP. Remarkably, summing over many NW modes yields $I_{wire}$ which is almost symmetric in $I_{SQUID}$, in agreement with the experiment.

Figure 3c shows $I_{wire}^{A}$ versus $\Phi$ and $I_{SQUID}$ over a wider range of SQUID currents. Specifically, $I_{wire}^{A}$ turns out to be a non-monotonic function of $I_{SQUID}$, initially increasing then being suppressed for large $I_{SQUID}$. This is emphasized in Fig. 3d where $I_{wire}^{A}(\Phi)$ is plotted for selected values of $I_{SQUID}$. $I_{wire}^{A}$ is a sinusoidal-like function of $\Phi$ whose amplitude depends on $I_{SQUID}$, and is maximized at $\Phi \sim (1/4)\Phi_{0}$ and $\Phi \sim (3/4)\Phi_{0}$.

The full $I_{wire}^{A}(V_{SQUID})$ dependence for a few values of flux is displayed in Fig. 3e and highlights both the monotonic linear increase for low $V_{SQUID}$ and suppression at large $V_{SQUID}$. The symmetry in $V_{SQUID}$ (i.e., in $I_{SQUID}$) is emphasized as well. Furthermore, $|I_{wire}^{A}|$ is maximized at $|V_{SQUID}^{\max}| \approx 0.4$ mV independently of $\Phi$, where it reaches values exceeding 20 pA. By converting $V_{SQUID}^{\max}$ in terms of the Josephson frequency we get $\nu_{J} \approx 190$ GHz whose corresponding time, $\tau_{D}^{-1} = 5$ ps, is comparable to $\tau_{D} = W^{2}/D \approx 4$ ps, i.e., the time required by electrons to diffuse in the NW between the Josephson junctions. In the above expression $W \approx 250$ nm is the width of the SQUID central electrode (Fig. 1b) which we assume to coincide with the separation between the weak-links, whereas $D \approx 0.015$ m$^{2}$/s is the diffusion coefficient of the NW [26]. The transition between the regime of $I_{wire}^{A}$ enhancement as a function of $V_{SQUID}$ to the one of $I_{wire}^{A}$ suppression can be explained in terms of the ability of the electrons to follow adiabatically the time-dependent parameters up to a maximum frequency set by $\tau_{D}^{-1}$. Another possible contribution to the suppression observed at larger $V_{SQUID}$ might stem from weakening of the ac Josephson coupling at high applied current [27].

The $I_{wire}^{A}(V_{SQUID})$ dependence plotted over a reduced bias range is displayed in Fig. 3f. In particular, $I_{wire}^{A}$ shows a linear behavior with slope $\eta$ which depends on the applied flux, and obtains values as high as several $10^{-3}$ pA/GHz. In the so-called ‘adiabatic regime’, i.e., where pumped current is expected to vary linearly with frequency, $\eta$ would therefore correspond to some $10^{-3}$ electrons per pump cycle.
A wire

rent due to asymmetry between the junctions which is predicted for quantum pumping (see Supplementary Information). This might pave the way to the investigation of the interplay between superconductivity-induced quantum pumping and exotic electronic states existing, for instance, in graphene [28] or in carbon nanotubes [29].

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METHODS SUMMARY

Selenium doped InAs NWs were grown by chemical beam epitaxy on an InAs 111B substrate. Gold catalyst particles were formed by thermal dewetting (at 520 °C for 20 min) of a 0.5-nm-thick Au film under TBA flux. NWs were grown for two hours at 420 °C using TBA, TMI and DTBSe metalorganic precursors with line pressures of 2.0 Torr, 0.3 Torr, and 0.4 Torr, respectively. NWs have diameters of 90±10 nm and are around 2.5 µm long.

The role of temperature (T) is shown in Fig. 4a which displays $I_{\text{wire}}^A$ versus $V_{\text{SQUID}}$ at $\Phi = (3/4)\Phi_0$ for several increasing temperatures. $I_{\text{wire}}^A$ monotonically decreases upon increasing temperature, which can be ascribed to the influence of thermal smearing as well as thermal-induced dephasing, and is suppressed for $T \gtrsim 3.5$ K. We stress that the aforementioned temperature is substantially smaller than $T_c$, the latter setting the disappearance of both Josephson effect and superconductivity in the JQEP. The $I_{\text{wire}}^A(T)$ dependence at the same flux is shown in Fig. 4(b) for a few $V_{\text{SQUID}}$ values. Specifically, $I_{\text{wire}}^A$ begins to round off at lower temperatures indicating a saturation, whereas it is damped at higher $T$. Low-temperature behavior suggests that current tends to saturate upon reducing temperature when the “effective” separation between Josephson junctions becomes of the same order of the electron coherence length in the NW, $L_T = \sqrt{\hbar D/(2\pi k_B T)} \sim 270$ nm at 250 mK, where $h$ is the reduced Planck’s constant while $k_B$ is the Boltzmann’s constant. By contrast, the decay of $L_T$ at higher temperatures may be considered as one of the predominant decoherence mechanisms leading to $I_{\text{wire}}^A$ suppression. Further study is needed to clarify this point.

It is worthwhile to emphasize that other effects which might manifest in the JQEP would yield currents characterized by symmetries markedly different from the ones predicted for quantum pumping (see Supplementary Information). Among these we recall (1) any spurious current due to asymmetry between the junctions which is always dominated by a component symmetric in $\Phi$ and antisymmetric in $I_{\text{SQUID}}$; (2) any thermocurrent generated by a different power dissipated in the two junctions, which is expected to be predominantly symmetric in both $\Phi$ and $I_{\text{SQUID}}$.

We finally note that other normal conductors than InAs NWs could be used for the implementation of the JQEP. The devices were fabricated using a technique of dry cleavage of the NWs followed by electron-beam deposition of Ti/V (15/120 nm) in an UHV chamber [21]. InAs NWs were treated with a NH$_4$S$_2$ solution before each evaporation step to get transparent metal-NW contacts [26].

The magneto-electric characterization of the devices was performed in a filtered $^3$He refrigerator (two-stage RC- and $\pi$-filters) down to $\sim 250$ mK using a standard 4-wire technique. Current injection at the SQUID terminals was obtained by using a battery-powered floating source, whereas voltage and current were measured by room-temperature preamplifiers. Derivative measurements (flux-to-voltage as well as flux-to-current transfer functions) were performed with standard low-frequency lock-in technique by superimposing a small modulation to the applied magnetic field.
pumping parameters varying along a closed path in the hybrid systems \[23–25\]. If \(x\) through a generalization of the Brouwer’s formula \[3\] to \(i = 1\)

space, and

where \(\Omega\) is the area enclosed by the path in parameter space, and

\[
Q_p = \frac{e}{2} \int \Omega \prod_{i,j} \alpha_i(x_1, x_2),
\]

(1)

where \(\Omega\) is the area enclosed by the path in parameter space, and

\[
\prod_{i,j} \alpha_i(x_1, x_2) = \prod_{i,j} \frac{\partial [S_{\text{he}}^*]_{i,j} \partial [S_{\text{he}}]_{i,j}}{\partial x_1 \partial x_2},
\]

(2)

In Eq. (3), \(S_{\text{ee}}\) and \(S_{\text{he}}\) are, respectively, the normal and the Andreev scattering matrices between the two N leads evaluated at the Fermi energy. Assuming that all leads support a single propagating channel, \([S_{\text{ee}}]_{i,i}\) and \([S_{\text{he}}]_{i,i}\) is the amplitude for an electron entering from lead \(i = 1, 2\) to be reflected back as an electron (hole), while \([S_{\text{ee}}]_{j,i}\) and \([S_{\text{he}}]_{j,i}\), with \(j \neq i\), is the transmission amplitude for an electron entering from lead \(j\) as an electron (hole). \(S_{\text{ee}}\) and \(S_{\text{he}}\) can be determined by the scheme proposed in Ref. \[23\], which requires the calculation of the scattering matrix \(S\) of the system depicted in Fig. 5 when all contacts are in the normal state. \(S\), in turns, is computed as a composition of three (three-legged) beam splitters, indicated by dashed circles in Fig. 5 and labelled by the index \(i = \{t, m, b\}\), connected to each other through a pair of ballistic N wires of different lengths. The scattering matrix of beam splitter \(\lambda\) can be written as

\[
S_{\lambda} = \left( \begin{array}{cc}
\sqrt{\gamma S_{\lambda}} & \sqrt{\gamma S_{\lambda}} e^{i\alpha_i} \\
\sqrt{\gamma S_{\lambda}} e^{-i\beta_i} & \sqrt{\gamma S_{\lambda}}
\end{array} \right),
\]

(3)

where \(\gamma\) takes values between 0 and 1/2, \(\alpha_i = -\psi + q_i \cos[-\gamma_i/(1 - \gamma_i)]\) and \(\beta_i = -\psi - q_i \cos[-\gamma_i/(1 - \gamma_i)]\) with \(q_i = \pm 1\). The three S leads, described by constant pair potentials \(\Delta_i = |\Delta| \exp(i\psi_i)\) (with \(i = 3, 5\)), are assumed to be ideally coupled to the system so that perfect Andreev reflection occurs at the S interfaces. When the bias current \(I_{\text{SQUID}}\) is larger than the critical current of the SQUID, a voltage \(V_{\text{SQUID}}\) develops across the latter. For the SQUID we assume the RSJ voltage-current relation \[22\]

\[
V_{\text{SQUID}}(\delta \varphi) = \text{sign}(I_{\text{SQUID}}) \sqrt{I_{\text{SQUID}}^2 - I_c(\delta \varphi)^2},
\]

(4)

where \(R\) is the total shunting SQUID resistance and \(I_c(\delta \varphi)\) is the flux-dependent SQUID critical current. The latter, used to fit the data in Fig. 1c, can be written as

\[
I_c(\delta \varphi) = (I_{c1} + I_{c2}) \sqrt{\tau^2 + (1 - \tau^2) \cos^2(\delta \varphi/2)},
\]

(5)

where \(I_{c1}\) and \(I_{c2}\) are the critical currents of the individual Josephson junctions composing the SQUID, and \(r = (I_{c1} - I_{c2})/(I_{c1} + I_{c2})\) is the degree of asymmetry of the SQUID.

From a practical point of view, we first calculate \(Q_1\) and \(Q_2\) through Eq. (1) assuming that the N leads 1 and 2 and lead 5 are grounded, while S leads 3 and 4 are kept at the potential \(V_{\text{SQUID}}\). This choice sets the phases of the superconductors as follows:

\[
\phi_3 = \text{sign}(I_{\text{SQUID}}) \omega_3 t
\]

(6)
so that no net current flows between the two parts of the circuit. Since the N and S parts dependently of \( \delta\varphi \) that the two parameters are maximally out of phase, in-expanding periodicity. The flux-antisymmetric component of the pumped current is obtained as \( I_{\text{wire}} = (\omega_j/2\pi) \left[ Q_p(\delta\varphi) - Q_p(-\delta\varphi) \right]/2 \). In Fig. 3b \( I_{\text{wire}}^A \) is plotted in units of \( I_{\text{wire}}^\text{max} = I_{\text{c1}} + I_{\text{c2}} \) and \( R_K = 2\pi h/e^2 \) is the Klitzing resistance. The current has been computed assuming that the NW carries 50 independent channels, each of which described by a scattering matrix obtained taking \( \psi_\lambda, q_\lambda \) and the phases accumulated along the two N wires as random parameters, while setting \( \gamma_\lambda = 1/10, \gamma_m = 1/11 \) and \( \gamma_b = 1/13 \).

In the configuration where lead 1 is a voltage probe (rather than connected to ground) one can calculate the voltage \( V_p \) which develops at lead 1 as a consequence of the charge pumped. \( V_p \), determined by setting to zero the current flowing in the NW, can be written as

\[
V_p(\delta\varphi) = |V_{\text{SQUID}}(\delta\varphi)| \frac{G_1(\delta\varphi) + G_2(\delta\varphi)}{G_1(\delta\varphi) G_2(\delta\varphi)} Q_p(\delta\varphi).
\]

The flux-antisymmetric component of \( V_p \) is defined as \( V^A_{\text{wire}}(\delta\varphi) = [V_p(\delta\varphi) - V_p(-\delta\varphi)]/2 \).

We shall further discuss the spurious effects which can occur in the presence of a shunting dissipative current across the Josephson weak-links. If the two Josephson junctions are not equal, a spurious voltage \( V_s \) (containing a constant and a time-oscillating component) arises in the NW between the beam splitters \( t \) and \( b \) in Fig. 5. This produces a current \( I_s \) in the NW that is not originated by quantum pumping. On the one hand, the current \( I_{s,\text{const}} \) related to the constant component of \( V_s \) reverses by changing the sign of \( I_{\text{SQUID}} \), in contrast to \( I_{\text{wire}}^A \) and it is an even function of \( \delta\varphi \). On the other hand, it turns out that the quantum rectified current \( I_{s,\text{rect}} \) associated to the oscillating component of \( V_s \) has no definite parity both in \( \delta\varphi \) and \( I_{\text{SQUID}} \), similarly to \( Q_p \) of Eq. (9). However, \( I_{s,\text{rect}} \) exists only in the presence of a finite \( I_{s,\text{const}} \), since they have the same physical origin. Yet, \( I_{s,\text{rect}} \) is smaller than \( I_{s,\text{const}} \) because the amplitude of the oscillating components of \( V_s \) is set by \( I_{\text{wire}}^\text{max} \), whereas the constant component of \( V_s \) is proportional to \( V_{\text{SQUID}} \). Therefore, the total spurious current is dominated by the component that is even in flux and odd in \( I_{\text{SQUID}} \) which would be detected, if present, in the transfer function \( I_{\text{wire}}^\text{A} \). Since the measured derivative signal \( I_{\text{wire}}^\text{A} \) is almost flux-symmetric [see Fig. 2(c)], we can rule out the presence of \( I_{s,\text{rect}} \) and therefore of quantum rectification. We stress that even in the presence of a sizable \( I_{s,\text{rect}} \), our calculations predict \( I_{\text{wire}}^\text{A} \) to be typically several orders of magnitude larger than the flux-antisymmetric component of \( I_{s,\text{rect}} \) (which is even in \( I_{\text{SQUID}} \)) thus fully dominating the measured signal.

In analogy, the current \( I_{\text{SQUID}} \) might produce a different power dissipated between points \( t \) and \( b \) in Fig. 5 leading to a thermocurrent flowing through the NW. Since \( V_s \) is dominated by its constant component, this thermocurrent would be almost symmetric both in \( \delta\varphi \) and \( I_{\text{SQUID}} \), in contrast to \( I_{\text{wire}}^\text{A} \). In addition, there could be a small contribution to the thermocurrent due to the oscillating component of \( V_s \) which would have no definite parity both in \( I_{\text{SQUID}} \) and \( \delta\varphi \). Since the power dissipated is proportional to \( V_s^2 \), such contribution to the thermocurrent is a fortiori negligible.

In conclusions, all the mechanisms envisioned above to produce a spurious dc current can be distinguished from quantum pumping by their parity with respect to magnetic flux \( \Phi \) or bias current \( I_{\text{SQUID}} \).

**Supplementary data** Here we present additional data for another JQEP device with nominally-identical geometry. Its essential parameters are the SQUID normal-state resistance of \( \sim 187 \Omega \) and the resistance of the Au/NW/Au line of \( \sim 2.1 \) k\( \Omega \). The general behavior of this device is similar to that discussed in the main text although it is characterized by less symmetry between the two Josephson junctions. Figure 6(a) displays the...
Fig. 6. Experimental data for a different JQEP. (a) $\Phi$-dependent modulation of the SQUID critical current $I_c$. Dashed line is the theoretical behavior of a tunnel and resistively-shunted junction SQUID assuming an asymmetry $r \sim 9\%$ between the critical currents of the two weak-links. (b) Color plot of the NW flux-to-current transfer function $I_{wire} = \partial I_{wire}/\partial \Phi$ versus $\Phi$ and $I_{SQUID}$. (c) Color plot of $I_{wire}$ versus $I_{SQUID}$ and $\Phi$. (d) $I_{wire}$ versus $V_{SQUID}$ for a few selected values of $\Phi$. Data in (a)-(d) are taken at $T = 250$ mK. (e) $I_{wire}$ versus temperature $T$ at selected bias currents $I_{SQUID}$ for $\Phi = (3/4)\Phi_0$. The error bars represent the standard deviation of the current values calculated over several measurements, and dashed lines are guides to the eye.

The full $I_s(\Phi)$ dependence of the SQUID measured at 250 mK which shows a maximum critical current of $\sim 330$ nA. Superimposed for a comparison (dashed line) is the model for a tunnel and resistively-shunted junction SQUID [22] assuming an asymmetry $r \sim 9\%$ between the critical currents of the two weak-links. The low-temperature flux-to-current transfer function $I_{wire} = \partial I_{wire}/\partial \Phi$ versus $\Phi$ and $I_{SQUID}$ is shown in Fig. 6. $I_{wire}$ shows no definite parity both in $\Phi$ and $I_{SQUID}$ which stems from the presence of a spurious current $I_s$ in the NW, which might be attributed to the reduced symmetry of the SQUID junctions. Figure 6 shows the extracted $I_{wire}$ versus $\Phi$ and $I_{SQUID}$ at 250 mK which highlights both the non-monotonic dependence and symmetry in $I_{SQUID}$. The full $I_{wire}(V_{SQUID})$ dependence for a few selected values of $\Phi$ at 250 mK is displayed in Fig. 3, and emphasizes the overall symmetry in $V_{SQUID}$. For the present device $I_{wire}$ is maximized at $V_{SQUID} \approx 0.25$ mV where it obtains values exceeding $\sim 27 \mu$A. $V_{SQUID}$ corresponds to a Josephson frequency $\nu_J \approx 120$ GHz (and related time $\nu_J^{-1} \approx 8$ ps). This difference from the device presented in the main text could originate from a slightly larger width $W$ of the SQUID central electrode combined with a reduced NW diffusion constant which lead to an increased diffusion time $\tau_D$. The $I_{wire}(T)$ dependence at $\Phi = (3/4)\Phi_0$ is shown in Fig. 6 for a few selected $I_{SQUID}$ currents. Specifically, $I_{wire}$ is rounded off at low temperature, whereas it is strongly damped and suppressed for $T \gtrsim 3$ K. The general behavior of $I_{wire}$ and the arguments of the previous section therefore suggest that in this sample $I_{wire}$ is fully dominated by quantum pumping, although a small component of quantum rectification might perhaps be present as well.

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f.giazotto@ens-lyon.fr

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