





Microwave spectroscopy of Andreev states in InAs nanowire-based hybrid junctions using a flip-chip layout

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Josephson junctions based on semiconductor nanowires are potential building blocks for electrically tunable qubit structures, e.g., the gatemon or the Andreev qubit. However, an actual realization requires the thorough investigation of the intrinsic excitation spectrum. Here, we demonstrate the fabrication of low-loss superconducting microwave circuits that combine high quality factors with a well-controlled gate architecture by utilizing a flip-chip approach. This platform is then used to perform single-tone and two-tone experiments on Andreev states in in-situ grown InAs/Al core/half-shell nanowires with shadow mask defined Josephson junctions. In gate-controlled and flux-biased spectroscopic measurements we find clear signatures of single quasiparticle as well as quasiparticle pair transitions between discrete Andreev bound states mediated by photon-absorption. Our experimental findings are supported by simulations that show that the junction resides in the intermediate channel length regime.

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Over the course of the last decade, Josephson junctions based on superconductor-insulator/normal conductor-superconductor stacks have been used as building blocks for various kinds of quantum bit structures¹ such as the fluxonium, the transmon, or the gatemon as its electrically-tunable counterpart. All these structures, including the gatemon, have in common that they rely on the collective excitation of an extended circuit structure and only utilize macroscopic properties such as the supercurrent to manipulate and interact with the system.

In contrast to conventional Josephson circuits, there are mesoscopic systems, which have been already realized by means of different superconductor hybrid structures, including break junctions, carbon nanotubes, and semiconductor nanowires. Here, the characteristics of the system can be determined by a small number of well-resolved microscopic excitations that are spatially localized within the normal conducting section, so-called Andreev bound states^{2,3}. Transitions between a pair of such states can be utilized in two special types of quantum bits, i.e., the Andreev level qubit^{4–7} and the Andreev spin qubit^{8–12}. In addition study of the energy-phase dispersion of Andreev bound states can provide insight into details of the superconducting state of the system. This is of interest as the range of hybrid systems being studied is expanding from metal and semiconductor-based systems to new topological insulator or two-dimensional systems, partly motivated by proposals for novel superconducting states such as proximity induced topological superconductivity^{13,14}.

The successful realization of any system based on Andreev-bound states requires a deep and detailed knowledge of the energy-phase dispersion of the Andreev levels. In the past, tunneling measurements^{15,16} or Josephson junction spectrometers¹⁷ have been employed for the spectroscopic analysis of single Andreev-bound states. As an alternative technique, the Andreev bound state spectrum can be detected through coupling of a phase-biased junction to a microwave cavity, e.g., a coplanar waveguide (CPW) resonator. This two-tone spectroscopy may provide evidence of the properties of sought after exotic phenomenon such as topologically protected Majorana modes^{13,18,19}. The coupling of the device with the cavity can be readout using the dispersive shift of the resonance single-tone response and transitions between Andreev bound states can be mapped using two-tone spectroscopy techniques typically employed in circuit quantum electrodynamics^{6,20–22}.

Each new hybrid system has its own set of material processing requirements, e.g., substrate limitations, the sensitivity to chemicals and etchants, or the thermal budget. In particular, semiconductor nanowire-based devices can require special etching steps and the deposition of high- k gate dielectrics that are known to be counterproductive to the performance of superconducting microwave resonators. Additionally, state-of-the-art fabrication techniques of superconducting microwave circuits such as deep trench etching are usually highly detrimental for semiconductor devices²³. Here, we propose a flip-chip method, with the aim of separating the fabrication of the Josephson junction circuit and the microwave circuit. Regarding the latter, the key is the reduction of parasitic two-level systems by optimizing interfaces and surfaces^{23–27} and limiting the effective area on the microwave circuit that is exposed to the dielectric. In addition, one can select low-loss substrate materials such as sapphire to improve the intrinsic quality of the superconducting layer without the risk of a degradation or destruction of the nanowire Josephson junction due to electrostatic discharge.

The flip-chip approach reported here creates a platform for the coupling of superconducting circuits containing high Q-factor coplanar waveguide resonators to various hybrid devices, including selectively-grown structures such as topological insulators²⁸, which can then be used to perform spectroscopy of Andreev states in any highly transparent system. In order to demonstrate the feasibility and performance of our flip-chip approach, we investigated the coupling of the resonator to the device chip. Furthermore, we present single-tone, as well as two-tone spectroscopic measurements of an epitaxially grown InAs/Al half shell nanowire with a shadow mask, defined Josephson junction^{29,30}. We note that a similar flip-chip method has recently been reported for Josephson junction detector-based spectrometers³¹.

Results and discussion

Flip-chip technique. A schematic representation of our system, consisting of the microwave circuit on sapphire (resonator chip) and the Josephson junction device on high-resistivity silicon (100) (device chip), is shown in Fig. 1a. Here, a rf-superconducting quantum interference device (rf-SQUID), comprised of a gate-controlled nanowire Josephson junction in a superconducting

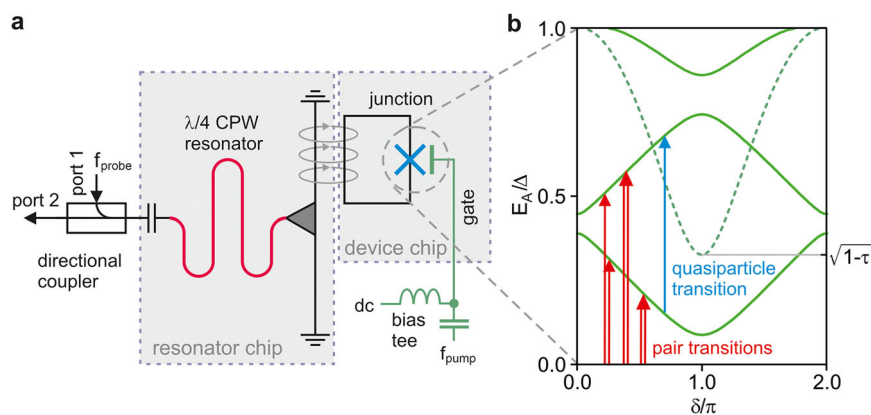


Fig. 1 Schematic sketch of the measurement setup and overview of the different excitations and transitions. **a** The nanowire Josephson junction on the device chip is inductively coupled to the $\lambda/4$ coplanar waveguide (CPW) resonator on the sapphire-based superconducting microwave circuit. The gate is connected to an external bias tee which supplies the dc-bias and the rf signal of the second tone. **b** Excitation representation of Andreev levels for a short junction with $L/\xi < 1$ (dashed line) and a long junction with $L/\xi = 2.74$ (solid lines). The transparency was set to typical state-of-the-art values for epitaxial nanowire Josephson junctions ($\tau = 0.895$). Here, L corresponds to the effective junction length, $\xi = \hbar v_F/\Delta$ is the superconducting coherence length, v_F the Fermi velocity, and Δ the superconducting gap. The blue and red arrows indicate possible single quasiparticle (odd) and quasiparticle pair (even) transitions, respectively.

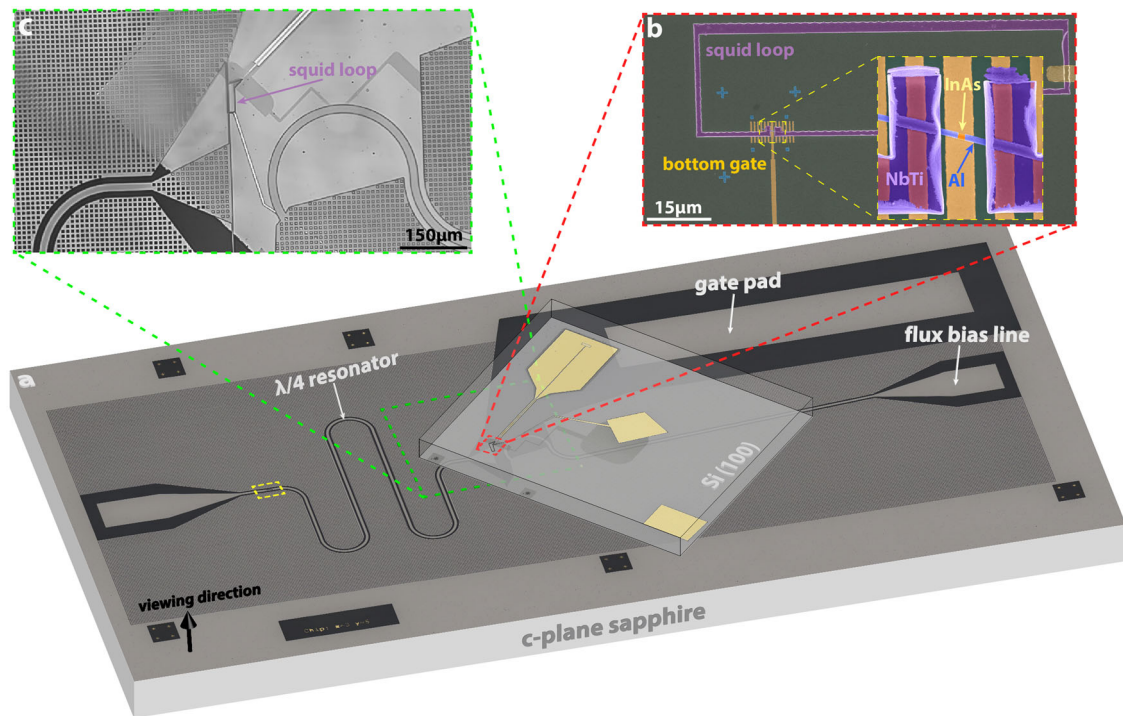


Fig. 2 Overview of an exemplary Flip-Chip device. **a** Optical micrograph of an exemplary microwave circuit with a 650 μm -thick sapphire wafer substrate in c-plane orientation. All superconducting structures are made out of 50 nm-thick molecular beam epitaxy grown Nb, which has been patterned by means of dry etching. For the application of the microwave signal, a single $\lambda/4$ coplanar waveguide resonator is used, which is capacitively coupled (yellow square) to a short transmission line and an impedance-matched superconducting bonding pad. The electromagnetic tuning of the device is realized with a combination of an on-chip flux bias line and a large bonding pad that acts as the mechanical and electrical connection between the gate electrode on the device chip and the fridge wiring. **b** Scanning electron microscopy image of an in-situ formed InAs/Al nanowire Josephson junction on top of a single bottom gate electrode. The SQUID loop itself and the electrical contacts are made out of 80 nm-thick NbTi. **c** Optical micrograph of a flipped device. The image was taken through the sapphire substrate mounted upon a polymer stamp during alignment. The SQUID loop is aligned such that one side is located above the termination point of the resonator. To ensure a constant and homogeneous reference potential, the device is connected to a large bonding pad that is electrically coupled to the ground plane of the sapphire chip with a conducting epoxy.

loop, is inductively coupled to a quarter wavelength ($\lambda/4$) CPW resonator. By this approach, it is possible to make use of well-defined pre-patterned surface gates and high-k dielectrics with low defect density on the device chip and still preserve high quality microwave structures. The superconducting Nb resonator is connected to a directional coupler, thus providing the pump tone as the input signal and feeding the reflected signal back to the measurement setup. The mesoscopic nanowire Josephson junction is based on an epitaxial Al-InAs half-shell nanowire Josephson junction device that has been prepared by utilizing an in-situ shadow evaporation process²⁹. In Fig. 2 an optical micrograph of the microwave circuit together with details of the alignment of the device chip on the resonator circuit as well as a micrograph of the nanowire-based Josephson junction are depicted. Further details about the mounting procedure are provided in Supplementary Methods 2.

When coupled to the microwave cavity, the resonator photons drive transitions between Andreev levels, and with sufficient coupling the state of the system is detected via the dispersive shift of the cavity resonance²⁰, i.e., the single-tone response. Here, two types of excitations of the Andreev system are possible. First, the microwave photon-assisted excitation of a pair from the superconducting condensate into the Andreev doublet with total energy ($E_{A,1} + E_{A,2}$), where $E_{A,n}$ is the energy of Andreev state n . Second, the excitation of a single quasiparticle from one Andreev level to another^{10,21,22}. Both transitions are in literature typically identified by the parity of the number of quasi-particles involved, i.e., even (pair) or odd (single). An example Andreev bound state

spectrum in the excitation representation is shown in Fig. 1b where each line indicates a doublet of spin degenerate states. The details of the energy dispersion depend not just on the superconducting gap and junction length but also properties of the weak link system such as spin-orbit interaction or confinement potential. For longer channels with $L/\xi > 1$, higher Andreev bound states are folded into the phase space, as indicated by solid lines Fig. 1b. The dashed line in Fig. 1b corresponds to the single channel/short junction limit, which has been utilized already to study Rabi oscillations in both atomic break junctions²⁰ and semiconductor nanowire-based junctions⁶. In the limit of a short junction, with length L much smaller than the superconducting coherence length ξ , a doublet of Andreev bound states appear at energies $E_A = \Delta\sqrt{1 - \tau\sin^2(\varphi/2)}$, where τ is the junction transparency and φ the superconducting phase difference^{32–34}.

Determination of resonator properties. As a first step to characterize the flip-chip set-up, the properties of the $\lambda/4$ CPW resonators are determined. Figure 3a provides an example for the single-tone response S_{21} of a highly overcoupled ($Q_c \approx 20,000$) resonator similar to the one shown in Fig. 2 with the nanowire device gated such that all states are detuned from the cavity and measured in a dilution refrigerator at $T \sim 15$ mK. For an applied microwave power of -80 dBm, a narrow dip in magnitude (solid blue line) can be observed at resonance. In addition, the resonator exhibits a pronounced phase shift at this frequency (dashed red line). Here, the phase signal has been corrected for the

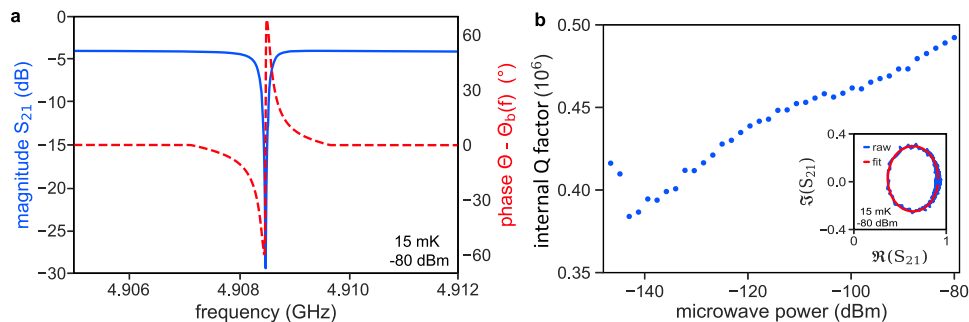


Fig. 3 Exemplary resonator response and power-dependence of the internal quality factor. **a** Single-tone response of a strongly overcoupled resonator similar to the microwave circuit shown in **a** with a sample chip attached to it and with the nanowire Josephson junction set into pinch-off. The blue line corresponds to the measured change in magnitude and the red line to the unwrapped phase, respectively. The latter has been corrected by the frequency-dependent background phase $\Theta_b(f)$. **b** Extracted internal quality factor of the very same resonator as a function of the applied microwave power. Measurements are performed at $T = 15$ mK and the rf-power is indicated with the fridge-internal attenuation accounted for. The inset shows an exemplary circle fit of the resonator in the IQ plane at an applied microwave power of -80 dBm.

monotonically changing, frequency-dependent background phase Θ_b . Figure 3b shows the power-dependent internal quality factor Q_i of the very same chip as in Fig. 3a. The microwave power is estimated at the resonator chip including 62 dB of cold attenuation inside the dilution fridge and external attenuation. The chip reveals a maximum of $Q_i = 5 \times 10^5$ in transmission-mode S_{21} . Based on an estimated single photon limit of ~ 150 dBm, our resonator reveals a large internal Q factor even at low photon numbers with the decrease in Q_i anticipated to arise from the increasing influence of unsaturated two-level systems in the device. The characteristic parameters of the resonator are thereby determined by means of the circle fit method³⁵ (see inset in Fig. 3b and Supplementary Methods 1 for further details). The device loaded Q is $Q_l \sim 26,000$. The internal Q-factor achieved for our cavity is moderate but with relatively simple one step dry etching fabrication. Increases in Q_i maybe realized through improved material processing or application of techniques such as deep reactive ion etching of the resonator channel²³ or superconducting ground plane air-bridges³⁶. The flip-chip approach has the benefit that such fabrication may be applied without concern for impact upon the processing of the Josephson junction.

Interaction of cavity and short channel junction. Before we present data of spectroscopic measurements it is useful to take a closer look at the predicted dispersive response of the cavity to a single Andreev bound states in the short channel case (cf. dashed line in Fig. 1b)^{4,6,21}. Figure 4a shows the energies for two-particle transitions with two very slightly different values of transparency τ . Shown in Fig. 4b, c are the corresponding predicted dispersive shifts of the cavity frequency in the weak coupling regime. The response is modeled applying a Jaynes-Cummings type Hamiltonian to an effective two-level system comprised of the ground state and the two quasi-particle excited state of the system, described by Zazunov et al.⁴ for which further details can be found in the Supplementary Note 2. As can be seen in Fig. 4b the cavity response is highly sensitive to the detuning of the two-particle transition to the Andreev bound state from the cavity frequency, $hf_A(\varphi) - hf_0$. When this detuning is small but positive the cavity resonance is dispersively shifted down in frequency. In contrast, if the detuning can be shifted from positive to negative through control of φ the resulting dispersive shift of the cavity shifts from downward to upward, as shown in Fig. 4c. Thus, from the type of cavity response, information on the detuning can be gained. Note that in the strong coupling regime, where the effective coupling g_{eff} is much greater than the lifetime of photons

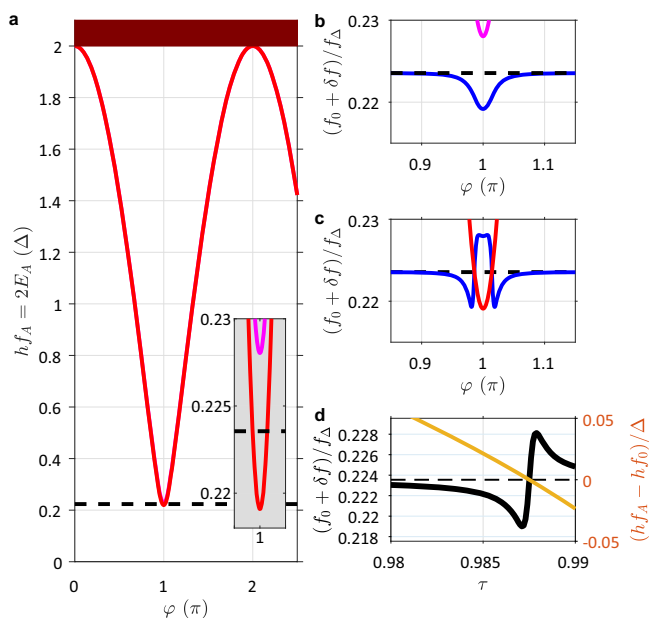


Fig. 4 Dispersive coupling of the cavity to a short channel and overview of the resulting frequency shift and detuning for different transparencies. **a** Pair transition energies (hf_A) for the short channel limit for two different transparencies. Magenta and dark purple traces indicate the quasi-particle transition energy hf_A for transparencies of $\tau = 0.987$ and 0.988 , respectively. **b, c** Cavity frequency shift δf from the bare cavity frequency f_0 normalized by the frequency corresponding to the gap edge $f_\Delta = \Delta/h$ for both $\tau = 0.987$ and 0.988 , respectively. The effective coupling strength at $\varphi = \pi$ is $g_{\text{eff}} = 0.0139f_\Delta$ for both transparencies. As can be seen in **b**, when the bound states and cavity detuning remains positive the cavity frequency is dispersively shifted downward. If however the detuning is shifted through zero via the phase as in **c** the cavity frequency exhibits both up and down shifts. Model parameters for **b** and **c** are: $\omega_c = 0.2235f_\Delta$, $\gamma = 0.0067f_\Delta$, $\kappa = 6.7 \times 10^{-6}f_\Delta$, $Z = 2 \times 10^{-5}$. Model details can be found in the methods section. **d** Cavity frequency shift δf as a function of τ at $\varphi = \pi$. Simulation parameters are $\omega_c = 0.2235f_\Delta$, $\gamma = 0.0067f_\Delta$, $\kappa = 6.7 \times 10^{-6}f_\Delta$, $Z = 2 \times 10^{-5}$. The detuning of Andreev bound state transition and cavity ($hf_A - hf_0$) is also plotted to highlight the point of reversal of frequency shift.

and Andreev states, such avoided crossing regions would manifest zero-mode Rabi splittings as has been previously reported in both atomic²⁰ and nanowire⁶ Josephson junctions. In Fig. 4d the relative frequency shift is plotted as a function of τ at $\varphi = \pi$

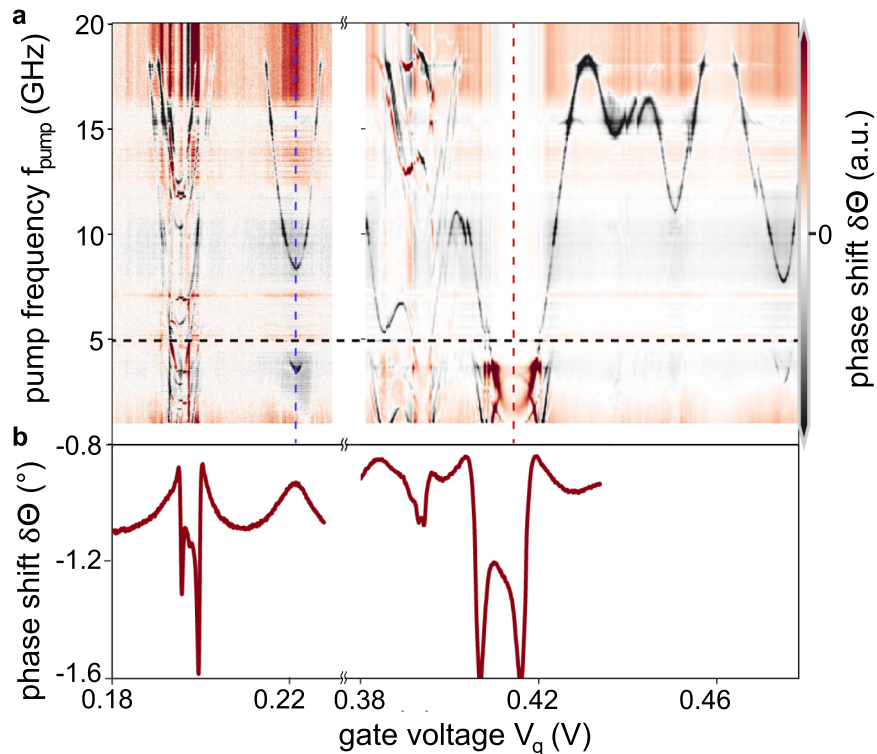


Fig. 5 Gate-dependent Andreev bound state spectrum. **a** Two-tone spectrum at $\varphi = \pi$ in a gate voltage range from 0.18 to 0.23 V and from 0.37 to 0.48 V. The phase shift of the resonator $\delta\Theta$ is color coded. The probe frequency is indicated by a dashed black line. We assign the abrupt discontinuities in the state dispersion, e.g., around $V_g = 0.43$ V, to random reconfigurations of the nanowire surface potential or changes in the effective gate charge. The vertical dashed lines indicate the gate voltages chosen for the flux-dependent measurements discussed later. **b** Corresponding resonator phase-shift obtained by single-tone spectroscopy in the same gate voltage range. The slight shifts compared to the measurement in **a** are caused by a time-dependent gate drift.

highlighting the strong dependence of frequency shift on the detuning.

Spectroscopic measurements. After discussing the properties of the $\lambda/4$ resonator coupled to the device chip, we now proceed with spectroscopic measurements on our InAs-nanowire/Al junction. As a first step, gate-dependent experiments were performed at a superconducting phase difference across the junction $\varphi = \pi$ to identify regions with transitions between Andreev levels. Figure 5a shows a two-tone spectrum as a function of gate voltage alongside a measurement of the phase shift in the single-tone response with only probe tone applied at the bare cavity frequency f_0 (cf. Fig. 5b). With probe tone at f_0 the phase shift scales linearly with the dispersive shift of the cavity frequency as long as that shift is small compared to the resonance width. The scaling is given by the loaded Q of the cavity with higher Q giving larger measured phase shifts for a given frequency shift. Note that due to the nature of the $\lambda/4$ cavity a negative phase shift indicates a positive frequency shift and vice versa.

The phase shift of the single-tone response (shown in Fig. 5b) reflects the detuning of the Andreev transitions from the cavity frequency. It reveals dispersive shifts of the cavity resonance due to interaction with the Andreev level(s), which are consistent with the states observed in the two-tone spectrum. The observed modulation of the single-tone response as a function of the applied gate voltage is caused by changes in the transition energy between different Andreev levels. At gate voltages for which the Andreev level transition energy is higher than the resonator photon energy (such as $V_g \sim 0.22$ V), the cavity resonance frequency is dispersively shifted down, giving a relative increase in phase. In contrast, at gate voltages for which the detuning

between Andreev transition energy and cavity frequency is reduced to zero, an abrupt reversal of the frequency shift is observed, inducing strong negative phase shifts (such as at $V_g \sim 0.41$ V). Around zero detuning, despite the signature of an avoided-crossing-like transition, we observe no zero-mode or vacuum Rabi splitting that would indicate strong coupling. The most likely reason are the comparably short lifetimes of the Andreev states being studied, as such the device is being probed in the weak coupling dispersive regime.

The two-tone spectra in Fig. 5a contains richer features than suggested in the single-tone response. For a careful selection of a gate voltage such as the region around $V_g \sim 0.195$ V or 0.39 V the two-tone spectrum reveals transitions to multiple Andreev bound states as well as multiple-photon processes. The multiple photon replicas of the quasiparticle pair transitions to the Andreev states are particularly evident in the low second tone frequencies below the cavity frequency (<5 GHz).

Proceeding to the lower gate voltage section in Fig. 5a one finds a rather complex spectrum with an abundance of transitions between $V_g = 0.19$ and 0.20 V. In addition, on its right side one observes a single spectral trace which is attributed to a pair transition. Once again, the coupling of the Andreev bound states in the junction to the resonator is reflected in the single-tone measurement shown in Fig. 5b. A strong modulation of the phase shift is observed at a gate voltage range from $V_g = 0.19$ to 0.20 V, where the transitions found in Fig. 5a cross the resonator frequency. In addition, around $V_g = 0.22$ V where the Andreev level approaches the resonator frequency the cavity phase shift has a smooth maximum.

One finds that at larger gate biases in the presented measurement range ($V_g = 0.42$ –0.47 V) only a single transition

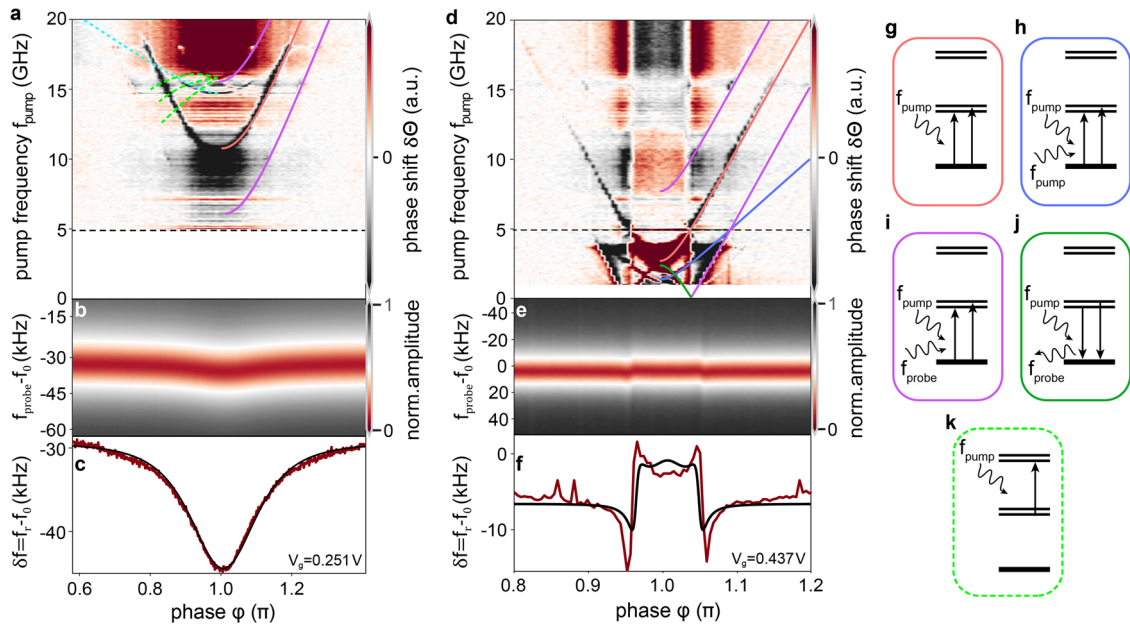


Fig. 6 Flux-dependent state dispersion for different gate voltages and schematic overview of the observed transitions. **a** Two-tone Andreev bound state spectrum at a gate voltage of $V_g = 0.251$ V (blue dashed line in Fig. 4 a) as a function of the magnetic field induced phase difference φ across the junction. The red line is a fit to the pair transition driven by the pump tone using the short channel model with parameters $\tau = 0.96 \pm 0.03$, $\Delta = 105 \pm 37$ μeV . The purple lines correspond to a corresponding two photon process caused by the absorption or emission of an additional resonator photon. The dashed blue line indicates another two-particle transition feature from another set of Andreev states. Dashed green lines highlight a set of single-quasiparticle transitions between spin polarized Andreev levels. **b** Single-tone measurement of cavity response at the same V_g as in **(a)**. The resonance frequency is extracted and plotted in **c** alongside a fit to a Jaynes-Cummings Hamiltonian (solid black line). Fit parameters are $\gamma = 9.7 \pm 0.1$ GHz, $Z = 1.09 \pm 0.02 \times 10^{-6}$, $f_{\text{offset}} = -29$ kHz, $g_{\text{eff}}(\pi) = 78 \pm 1$ MHz, and $\kappa = 0.1964$ MHz. **d** Two-tone spectrum at a gate voltage of $V_g = 0.4374$ V (red dashed line in Fig. 4a). Again a prominent two particle transition is observed and fitted with a short channel model with $\tau = 0.998 \pm 0.001$, $\Delta = 132 \pm 40$ μeV (solid red line). Two photon transitions (solid purple lines) which utilize both cavity and pump photons are also observed. Additionally, a solid blue line indicates a two photon pair transition using two pump photons. The solid green line corresponds to a higher-order transition due to the combination of a pump-induced single resonator photon and the relaxation of an Andreev level. **e** Single-tone measurement of cavity response at the same V_g as in **(d)**. The resonance frequency is extracted and plotted in **f** alongside a fit to a Jaynes-Cummings Hamiltonian (solid black line). Fit parameters are $\gamma = 0.9 \pm 0.2$ GHz, $Z = 4.8 \pm 0.7 \times 10^{-8}$, $f_{\text{offset}} = -8.4$ kHz, $g_{\text{eff}}(\pi) = 22 \pm 3$ MHz, and $\kappa = 0.1964$ MHz. **g-k** Schematics of the different transitions being discussed in **a** and **d** showing quasi-particle excitations from Andreev levels (two parallel lines) or the ground state (solid line) and the pump or probe photons that contribute to the process.

is observed in the pump frequency window. We attribute the single transition found at gate biases $V_g > 0.42$ V to a quasiparticle pair transition $2E_A = hf_{\text{pump}}$ from the ground state^{21,37}. The modulation with the gate voltage can be attributed to potential fluctuations in the nanowire channel, e.g., in a chaotic open dot with low charging energy coupled to superconductors by tunnel barriers and a diffusive junction with transparent normal superconducting interfaces, as discussed by Houzet and Skvortsov³⁸. Indeed, in comparable nanowire-based junctions critical current fluctuations were observed, which were correlated to corresponding conductance fluctuations in the normal state^{39,40}. Since ultimately the critical current is determined by the Andreev levels a variation of these levels with gate voltage will directly result in critical current variations. While only a single pair transition is observed in the two-tone spectrum for $V_g > 0.42$ V we deduce that further high lying transitions must exist at frequencies beyond $f_{\text{pump}} = 20$ GHz. The spectroscopy as a function of flux (discussed below) reveals some characteristics typical of long Josephson junctions and relatively large frequency shifts from the bare cavity frequency across the measurement flux range are observed.

We next discuss the phase-dependent measurements and use a combination of single and two-tone measurements to evaluate the coupling of the Andreev levels. In Fig. 6a, d are shown for which we select gate voltages of $V_g = 0.251$ V and 0.4347 V, respectively. The spectrum at $V_g = 0.251$ V corresponds to a gate condition at

which the pair transition frequency remains above the cavity frequency at $\varphi = \pi$. In contrast at $V_g = 0.4347$ V the pair transition frequency dips below the cavity frequency at $\varphi = \pi$. Note that owing to a gate drift the gate voltage was slightly readjusted between phase-dependent and gate-dependent measurements. The most prominent feature in both sets of phase data is the pair transition where $2E_A = hf_{\text{pump}}$. The pair transition can be fitted using solutions to a transcendental (non-algebraic) expression for the Andreev level energy in a long junction with multiple transverse 1-dimensional normal bands, as discussed in refs. 21,37 (see Supplementary Note 3). Features can also be well fit using the short channel analytical solution for the Andreev state energies if the gap Δ is taken as a free parameter which is reduced below the expected bulk gap (assumed to be $\Delta_{\text{Al}} \sim 180$ μeV based on previous dc transport studies²⁹). We note both fitting methods produce comparable results in the measurement parameters space but differ at higher pump frequencies. In Fig. 6a, d fits are performed using the short channel expression shown as solid red lines with parameters as indicated in the caption and results are later used to evaluate the coupling strength assuming coupling with only a single two-level system. The reduced gap required for fitting is a consequence of the states themselves being intermediate between the short channel limit and a long junction (see the Supplementary Note 1 for further discussion). In either case, we observe a high transparency $\tau = 0.995$ and 0.998 with fitted $\Delta \sim 104$ and 132 μeV in Fig. 6a, b respectively. We note that

fitting with the long junction expression also produces comparable high transparencies (see Supplementary Note 1).

Apart from the pair transition features in Fig. 6a, d, we also find weaker transitions. First, we identify replicas of second order transitions involving absorption or emission of a resonator photon, i.e., at $2E_A \pm f_{\text{res}} = hf_{\text{pump}}$, respectively (solid magenta lines). Furthermore, one finds a transition involving a pair of pump photons, i.e., $2E_A = 2hf_{\text{pump}}$ (dashed blue line). The appearance of these replicas is likely related to relatively high power of both pump and probe tones. Interestingly, close to $\varphi = \pi$ we also find a transition with an inverted dispersion (solid green line). We attribute this feature to a higher-order transition where a single resonator photon is created upon pump irradiation with the assistance of the Andreev energy. The relatively prominence of this feature also suggests elevated thermal occupation of the excited states of the Andreev system.

In addition to the two-tone measurement we present the single-tone response in Fig. 6b and e, as well as the extracted positions of the resonance frequency, plotted as shift $\delta f = f_{\text{res}} - f_0$ in Fig. 6c, f. Here, the resonances are fitted with a simple Lorentzian lineshape with a linear background subtraction to extract f_{res} . Note that a negative frequency shift seen here is consistent with a positive phase shift in data plotted in the previous gate spectra (cf. Fig. 5) owing to the features of the resonator. It can be seen that the direction of the resonance frequency shift is related to the detuning of the pair transition frequency and bare cavity frequency, as previously discussed. Most prominent modulation of the resonance occurs when detuning is tuned through zero, as seen in Fig. 6f, where the resonance frequency shift changes sign (cf. Fig. 4c). We fit the extracted frequency shifts assuming that only the single pair transition identified in the two-tone measurement contributes to the cavity response using the known analytical expressions for the coupling in the short channel limit^{4,20}. This approach is arguably an oversimplification of our system as other states are inevitably involved as will be later discussed but nevertheless serves as an indication of the lower limit of the coupling and a point of comparison with other studies. Fit use the extracted τ and Δ from two tone fits and the cavity photon lifetime as fixed parameters. Fits to Fig. 6c indicate an effective coupling strength at $\varphi = \pi$ of $g_{\text{eff}}(\pi) = 78 \pm 1$ MHz and a two-level system lifetime $\gamma \sim 9.7 \pm 0.1$ GHz. Fits to Fig. 6f indicate $g_{\text{eff}}(\pi) = 22 \pm 3$ MHz and a two-level system lifetime $\gamma \sim 0.9 \pm 0.2$ GHz. We note that fitting includes an offset to the frequency shift independent of φ that is -29 and -8.4 kHz in Fig. 6e, f, respectively. This shift is likely already caused by the coupling with additional states that are possibly beyond the range of applied pump tone.

Close inspection of the two-tone spectrum in Fig. 6 reveals a set of lines (dashed green) around 15 GHz which can be assigned to single quasiparticle transitions^{10,21,37}. These four-fold transitions originate from modifying the dispersion of the Andreev bound states by Rashba spin-orbit coupling. More precisely, the coupling of lowest Rashba-split confinement state in the nanowire to next higher ones results in a difference of Fermi velocities for different spin orientations. This in turn results in an effective spin-splitting of the Andreev bound states. Single quasiparticle transitions between pairs of these Andreev bound states then lead to a four-fold transition pattern in the two-tone spectroscopy, as found in our experiment. Single particle transitions between two subbands of the Andreev spectrum requires a relatively long junction and is thus inconsistent with the pair transition states previously discussed. In addition a weak pair transition like feature with a flatter dispersion (dashed blue line) is observed, again consistent with a junction that is effectively longer. It is found to be challenging to produce accurate fits to the

experimentally observed features using expressions for the long junction detailed by Tosi et al.^{21,37} which assumes two bands with different Fermi velocities due to the influence of spin-orbit interaction characterized by parameter $\Lambda_{j=1,2} = L\Delta/hv_{F,j}$ for band j . A rough qualitative matching is obtained by assuming that the single quasiparticle transitions are in-fact a two photon process involving a cavity and pump photon and an example of such a spectrum is shown in Supplementary Note 3. Model parameters include $\tau = 0.64$ and $\Lambda_1 = 3.17$, $\Lambda_2 = 2.06$ with $\Delta = 180 \mu\text{eV}$. We conclude that these feature arise from a lower transparency and relatively longer channel for which such single particle transitions may become visible owing to the reduced separation in energy between the Andreev subbands. We propose two origins for the coexistence of such long and short junction behavior. It is possible that states are related to two different transverse bands of the nanowire with a significant difference in the Fermi velocity. Alternatively, the states may correspond to channels at different locations on the nanowire surface. The nanowire is half covered with the Al shell such that transport on the nanowire surface that is on the underside (the opposite face from regions with aluminum deposition) may have an effective longer junction length. However, we would like to point out that the observed features differ from the ones observed in ref. ²¹ as only three lines are visible. Recent studies of ref. ²² indicate that this could be caused by additional junction-internal effects such as electron-electron interaction and thus highlight this behavior as a target for future studies.

Conclusion

We have demonstrated a flip-chip method for the coupling of a rf-SQUID device with a high-Q coplanar waveguide microwave cavity for spectroscopic measurements. The separate processing of the rf-SQUID device chip and resonator chip simplifies fabrication and relaxes requirements, i.e., for the resonator chip the low-loss substrate material could be chosen to improve the intrinsic quality of the superconducting resonator without taking the risk of degradation or destruction of the Josephson junction on the device chip. Using a demonstrator flip-chip device containing an Al-InAs half-shell nanowire Josephson junction, we were able to show that detailed information about Andreev bound states can be obtained with our approach using the single- and two-tone response. We find effective couplings of single Andreev-bound states to the cavity that are comparable to other reports on non flip-chip circuits. The concept presented here, can also be employed for more delicate or complex device circuits, e.g., ones containing topological insulator films, where the thermal budget is rather limited, or circuits with a set of gatemon or Andreev qubits.

Methods

Fabrication of the microwave circuit. The circuit containing the resonator is fabricated onto a *c*-plane sapphire substrate²⁷. The $\lambda/4$ resonator is designed in a coplanar waveguide (CPW) geometry with a bare resonance frequency of $f_0 \sim 5.4$ GHz. In order to avoid parasitic effects of residual resist or lift-off induced sidewalls, all superconducting structures have been fabricated out of a globally deposited 50 nm thick sputtered niobium film subsequently patterned with a SF₆ dry etching procedure. The chosen resonator design utilizes a single port and is arranged with meanders to occupy half of the available 8×4 mm² sapphire substrate, as shown in Fig. 2a. The remaining half of the substrate provides ample space on which to flip the target device chip while minimizing the overlap of device chip and resonator transmission line. In addition, space for gate connections and an additional microwave line for input of additional signals is provided. In order to stabilize induced vortices in an out-of-plane magnetic field, a global array of $4 \times 4 \mu\text{m}^2$ wide flux traps was etched into the ground plane and the termination point of the resonator. In Fig. 2c, a detail of the target device chip aligned to the termination point of the $\lambda/4$ coplanar waveguide resonator prepared on the sapphire substrate is shown.

Fabrication of the device chip and the nanowire Josephson junction. The target device chips are fabricated on a Si(100) substrate with high resistivity ($\rho > 100$ k Ω cm) to reduce substrate-induced losses. Local bottom gates are defined using Ti/Pt (3/10 nm) and are covered with a 3 nm/12 nm thick Al_2O_3 and HfO_2 dielectric layer that has been globally deposited by atomic layer deposition. Using a CHF_3/O_2 -based dry etching procedure, the latter is then removed from the areas that are not used for the actual formation or control of the SQUID loop or the nanowire Josephson junction. Single InAs/Al nanowire-based junctions are picked up and positioned on the back-gate using a micromanipulator mounted inside a scanning electron microscope (Hitachi FIB NB-5000). The junctions were fabricated by employing a fully in situ shadow evaporation scheme²⁹. The InAs nanowire has a diameter of 90 nm and a length of 6 μm . The 20 nm-thick Al half-shells are separated by 90 nm. After the nanowire placement, a superconducting loop of sputtered NbTi with a thickness of 80 nm is then fabricated using electron beam lithography and DC magnetron sputtering. In addition to the coupling loop, a ground connection pad is added to ensure a ground reference for the applied gate voltage. A scanning electron micrograph of the device circuit is shown in Fig. 2b.

Mounting procedure. The mounting of the target device chip upon the resonator is achieved using a homemade polymer stamping system. The sapphire resonator is mounted on a section of polymer gel sheet (Hakuto Gel Sheet WF-55-X4-A) stuck to a glass slide. Viewing through the back of the sapphire resonator the operator is able to align the device superconducting loop with the coupling section of the $\lambda/4$ resonator and bring the resonator and device chip into contact. The electrical connection between the two chips has been successfully achieved using two methods. In one a 2 μm -thick indium pad is added to the gate and ground contact regions of the sapphire chip and the device chip is heated to ~ 120 °C when brought into contact with the sapphire. Once the indium is observed to be displaced by the contact of the two chips the temperature is reduced. Alternatively, it was found that a small amount of silver epoxy (either Creative Materials 118-15 or Epo-tek E4110-PFC) at the gate and ground contact regions is sufficient to adhere the two chips and provide good electrical contact. Of the two methods, the indium contacts provide a stronger bond but the epoxy method is more forgiving in the event of a mistake in alignment, as this can be corrected before the epoxy is cured by carefully removing the epoxy, cleaning, and retrying.

Measurement setup. Spectroscopy measurements were performed in a BlueFors LD400 dilution refrigerator with a base temperature of 14 mK. The phase-bias (φ) across the junction was imposed by threading a small perpendicular magnetic field which is generated by a home-made superconducting coil through the superconducting NbTi loop. The magnetic flux is directly related to the junction phase difference through $\varphi = 2\pi\Phi/\Phi_0$, where Φ_0 is the magnetic flux quantum. The bottom-gate is connected to a bias-tee in order to apply a dc gate-bias V_g as well as to supply a microwave signal with frequency f_{pump} for two-tone spectroscopy. The device is characterized by single-tone as well as by two-tone spectroscopy. For the former the shift of the resonator frequency $f_r - f_0$ induced by the coupling to the weak link is measured by applying a continuous signal to the $\lambda/4$ resonator. The measured bare resonator frequency is $f_0 = 4.911019$ GHz. For two-tone spectroscopy, the pump tone is modulated using a fast microwave switch while the probe tone is continuous. The modulation of I and Q are readout using two lockin amplifiers with the pump modulation as a reference frequency. Additional details of the experimental setup can be found in Supplementary Methods 1.

Two-tone spectroscopy. Two-tone spectroscopy is similar to the concept of dispersive readout in terms of an indirect scanning of the coupled quantum system. A key difference is the way in which the second tone is applied to the device. Instead of the resonator, the impedance-matched gate line is used as a second and independent microwave channel. The measurement scheme is as follows: First, a frequency-dependent measurement of the cavity response is performed to find the resonance frequency f_0 of the resonator. The latter is then used as the first tone (probe signal) $f_{\text{probe}} = f_0$. In order to determine a range of parameters in which the resonator couples to one or more Andreev-bound states, the single-tone response is mapped out with respect to the applied gate voltage and magnetic field. Based on these results, a suitable gate voltage V_g or flux-induced phase difference φ is selected which is then kept constant throughout the measurement. Subsequently, the second tone (pump signal) f_{pump} is swept over a wide frequency range (here $\delta f_{\text{pump}} \approx \Delta_{\text{Al}}/4$). When the pump frequency is resonant with a transition between Andreev states the system will become excited through the absorption of a photon. This induces a significant alteration of the inductance of the SQUID loop and, in the following, induces a dispersive shift of the resonance frequency which may be detected as a change in phase $\delta\theta$ measured at the probe frequency. Lastly, the remaining electromagnetic variable, i.e., gate voltage V_g or magnetic flux Φ , is varied to change the energy of the bound state. The advantage of this approach is a very high sensitivity and the possibility to directly map out the energy space of the combined system far below and above the actual cavity resonance frequency. Assuming an undercoupled system, the only limitation is set by the bandwidth of the experimental setup and the wiring in the dilution refrigerator.

Fitting cavity frequency shifts. Fitting of the single-tone response of the cavity is performed using the modeled coupling of a short channel from prior studies^{4,20} in which the system is simplified to a Jaynes–Cummings type Hamiltonian. Applying input–output theory we find an expression for the cavity reflection:

$$A = \frac{-i\kappa}{(\omega_c - \omega_{\text{probe}}) - \frac{i\kappa}{2} + g_{\text{eff}}(\varphi)\chi} - 1,$$

where κ is the cavity loss including internal loss and the loss through port 1, $\kappa = \kappa_{\text{int}} + \kappa_1$. Term γ is the decoherence which can include both energy relaxation, poisoning and also pure dephasing rate. Term ω_{probe} is the probe frequency applied to the cavity and ω_c is the cavity resonance frequency. Parameter χ is the susceptibility given as

$$\chi = \frac{g_{\text{eff}}(\varphi)}{-((E_A/h) - \omega_c) + \frac{i\gamma}{2}}.$$

The effective coupling g_{eff} is a function of φ and is given as^{4,20}

$$g_{\text{eff}}(\varphi) = M \frac{1}{h} \sqrt{\frac{\hbar\omega_c}{2l}} \frac{\Delta E_A(\pi, \tau)}{4\Phi_0 E_A(\varphi, \tau)} \tau \sin(\varphi) \tan(\varphi/2),$$

where M is the mutual inductance between cavity and SQUID loop and l is the resonator length and $\Phi_0 = h/2e$. For fitting and simulation the expression is simplified to:

$$g_{\text{eff}}(\varphi) = \sqrt{Z} \frac{\Delta E_A(\pi, \tau)}{2 E_A(\varphi, \tau)} \tau \sin(\varphi) \tan(\varphi/2),$$

where $Z = (M^2/h^2)(\hbar\omega_c/2L)(1/2\Phi_0)$. Further details can be found in Supplementary Note 2.

Resonator characterization. In order to determine the characteristic properties of the resonators, namely the internal, coupling, and loaded quality factors Q_i , Q_c , and Q_l , the circle fit method presented in ref. 35 is applied to the resonator data. This approach employs the full scattering matrix of the VNA signal instead of only the magnitude or phase data and derives the system properties from the IQ plane, i.e., $\Re(S_{21})$ and $\Im(S_{21})$, using the fitted reflection coefficient Γ

$$\Gamma = \left[\frac{\Gamma_{\text{min}} + 2i \frac{Q_c Q_l Q_i}{Q_c Q_l + Q_l Q_i + Q_i Q_c} \frac{f - f_0}{f_0}}{1 + 2i \frac{Q_c Q_l Q_i}{Q_c Q_l + Q_l Q_i + Q_i Q_c} \frac{f - f_0}{f_0}} - 1 \right] e^{i\phi} + 1, \quad (1)$$

for an overcoupled resonator as a function of the probe frequency with

$$\Gamma_{\text{min}} = \frac{Q_l Q_c - Q_l Q_i + Q_c Q_i}{(Q_l + Q_i)(Q_c + Q_i)}. \quad (2)$$

Here, f refers to the applied probe frequency, f_0 is the resonance frequency and ϕ is a phase offset that describes the tilt of the circle in the IQ plane. The advantage of this non-iterative calculation is its independence of properly defined starting values. In addition, unwanted side effects such as the cable delay, phase noise, or the gain and attenuation along the fridge wiring are taken into account during the pre-processing of the scattering data.

Data availability

All the data that support the plots and the other findings of this study are available from the corresponding author upon reasonable request.

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Author contributions

P.P., B.B., and M.L.L. performed the nanowire growth by MBE and the deposition of the superconducting shell. P.Z. and R.S.D. fabricated the samples, carried out the low-temperature experiments and the subsequent data analysis. P.Z., R.S.D., and T.S. wrote the manuscript. All authors contributed to the discussions.

Competing interests

The authors declare no competing interests.

Additional information


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