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Engineering Strong Intrinsic Longitudinal Coupling in Circuit Quantum Electrodynamics



Robin Dekker

ENGINEERING STRONG INTRINSIC LONGITUDINAL COUPLING IN CIRCUIT QUANTUM ELECTRODYNAMICS

ENGINEERING STRONG INTRINSIC LONGITUDINAL COUPLING IN CIRCUIT QUANTUM ELECTRODYNAMICS

Proefschrift

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aan de Technische Universiteit Delft,
op gezag van de Rector Magnificus Prof. dr. ir. H. Bijl,
voorzitter van het College voor Promoties,
in het openbaar te verdedigen op woensdag 10 juni 2026 om 10:00 uur

door

Robin Christian DEKKER

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SUMMARY

Superconducting microwave circuits provide a versatile platform for engineering controllable quantum systems, with applications in quantum information processing, quantum sensing, and microwave optomechanics. In this thesis, we investigate how key circuit elements can be improved by engineering new interactions, enhancing microwave readout, and reducing dissipation in superconducting resonators. We achieve this through the design, fabrication, and cryogenic characterization of superconducting circuits, focusing on three themes: longitudinal coupling, parametric amplification for microwave measurements, and reducing microwave loss through surface processing.

In **Chapter 1** we introduce superconducting microwave circuits and motivate the main research directions of this thesis.

Chapter 2 describes the experimental workflow used throughout this thesis. Starting from circuit design and electromagnetic simulation, we explain how the devices are realized in the cleanroom, with particular attention to the practical details of the nanofabrication process. We then describe chip packaging, and give an overview of the cryogenic measurement setup used to characterize the devices.

Chapter 3 demonstrates how to implement an intrinsic longitudinal coupling between a transmon qubit and a linear microwave resonator by using the circuit quantum electrodynamics analog of radiation-pressure coupling. We discuss the experimental characterization of this interaction and highlight its prospects for high-connectivity quantum hardware and for experiments interfacing superconducting circuits with massive quantum systems.

Chapter 4 presents the design, fabrication, and characterization of a flux-tunable parametric amplifier based on the Superconducting Nonlinear Asymmetric Inductive eLement (SNAIL). We describe the device design and wafer-scale fabrication, and characterize the amplifier response across its tunable operating range. We investigate the gain-bandwidth trade-off and quantify the dynamic range through gain-compression measurements.

Chapter 5 investigates microwave loss in tantalum superconducting coplanar-waveguide resonators. We study surface-related dielectric loss in the single-photon regime using devices patterned from a single tantalum film with varying surface participation ratios. By comparing different resonator geometries and post-fabrication surface treatments, we show that additional cleaning steps can significantly reduce loss and yield reproducible improvements in the internal quality factor.

Finally, **Chapter 6** concludes this thesis by summarizing the main findings and providing an outlook for future research.

1

INTRODUCTION

Superconducting microwave circuits have emerged as a leading experimental platform for realizing and controlling quantum systems, with applications ranging from quantum sensing to quantum information processing. The foundations of quantum mechanics were laid in the early twentieth century, beginning with Planck's introduction of energy quantization [1]. In 1925, Heisenberg's formulation of matrix mechanics marked the start of modern quantum theory [2], which enabled quantitative descriptions of microscopic physical systems. Now, more than a century later, these ideas enable engineered quantum systems in which quantum behavior can be designed and tuned in controllable devices.

In superconducting quantum circuits, Josephson junctions provide a controllable nonlinearity that shapes the energy spectrum and interactions of microwave modes. This enables a broad range of devices, including superconducting qubits, quantum-limited parametric amplifiers, frequency converters, and tunable couplers. In the case of a superconducting qubit, the circuit is engineered to behave as an artificial atom, with its two lowest energy states forming a quantum bit (qubit). The first experimental demonstrations of superconducting qubits were realized in the late 1990s through coherent control of a Cooper-pair box [3]. This work showed that macroscopic superconducting circuits can exhibit coherent quantum dynamics and can be coherently controlled. Soon after, alternative implementations were explored, including the persistent-current (flux) qubit demonstrated here in Delft [4]. Together, these early experiments established that coherent quantum behavior can be realized across multiple superconducting circuit architectures.

Building on this progress, strong coupling between a superconducting qubit and a microwave resonator was achieved, enabling the coherent exchange of excitations between the two systems [5]. This result marked the start of circuit quantum electrodynamics (cQED), which adapts key concepts from cavity QED to superconducting circuits [6–8]. In cavity QED, strong coupling was experimentally identified through spectroscopic signatures such as vacuum Rabi splitting [9]. Demonstrating the same regime in superconducting circuits showed that engineered, on-chip systems could support the same fundamental light–matter interactions, while offering extensive control over system parameters.

The rapid development of superconducting circuit technology has since established cQED as a platform for exploring quantum optics in the microwave domain and for building scalable quantum processors [10–13]. Despite these successes, important conceptual and practical challenges remain. One longstanding conceptual problem is the incompatibility of quantum mechanics with general relativity. In quantum mechanics, time is treated as an external parameter that specifies the evolution of a system, whereas in general relativity, time is a dynamical component of spacetime that can depend on the gravitational field and differ between observers [14]. Motivated by this tension, our group investigates quantum behavior in increasingly massive systems, leveraging techniques from optomechanics and circuit quantum electrodynamics.

On the practical side, parasitic couplings complicate the design and control of multi-qubit systems, motivating alternative coupling schemes such as longitudinal coupling. Beyond improving quantum processors, longitudinal interactions can also be advantageous for optomechanical systems, where many protocols rely on strong coherent cavity drives and therefore require interactions that remain robust at large intracavity photon numbers [15–21]. Furthermore, the performance of superconducting quantum devices is often limited by loss and de-

coherence arising from coupling to the environment. In many cases, loss and decoherence are dominated by dielectric two-level systems (TLS) hosted in amorphous surface and interfacial layers, which can limit coherence times and resonator quality factors. Continued progress therefore depends not only on improving circuit design, but also on advances in materials processing and surface engineering.

This thesis addresses several of these challenges through the design, fabrication, and characterization of superconducting microwave circuits. It focuses on three main themes: (i) engineering an intrinsic longitudinal qubit–resonator interaction, (ii) parametric amplification, and (iii) reducing loss in superconducting resonators. Together, these topics span key aspects of superconducting circuits, from fundamental interactions to practical device improvements, and contribute to the continued development of superconducting quantum hardware.

1.1. THESIS OUTLINE

This thesis is organized as follows.

- **Chapter 2** describes the design, fabrication, and measurement methods used throughout this thesis. It introduces the device design workflow based on electromagnetic simulations and outlines the cleanroom processes used to fabricate superconducting devices. The chapter also describes the device packaging and cryogenic measurement setup.
- **Chapter 3** reports the experimental realization of intrinsic longitudinal coupling between a transmon qubit and a linear microwave resonator. The device concept and operating principle are introduced, and the emergence of longitudinal coupling from the circuit design is explained. Spectroscopic and time-domain measurements are used to activate and characterize a drive-induced Jaynes-Cummings interaction. The measurements demonstrate a longitudinal coupling rate that exceeds the relevant decay rates in the system, providing access to the strong coupling regime.
- **Chapter 4** presents the design, wafer-scale fabrication, and experimental characterization of a parametric amplifier based on the Superconducting Nonlinear Asymmetric Inductive eLement (SNAIL). We first investigate the linear (unpumped) properties of the device, followed by the demonstration of parametric amplification over a wide flux-tunable frequency range. The gain-bandwidth trade-off is examined, and the dynamic range is characterized through gain-compression measurements.
- **Chapter 5** focuses on a systematic study of microwave loss in tantalum coplanar waveguide resonators. We investigate how variations in the surface participation ratio and post-fabrication surface cleaning affect the internal quality factor in the single-photon regime. By comparing devices with different surface participation ratios and surface treatments, we identify the most effective cleaning procedures for reducing microwave loss, and show that this allows us to reliably obtain internal quality factors exceeding one million.
- **Chapter 6** summarizes the main results of this thesis and outlines the outlook for future research.

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2

DEVICE DESIGN, FABRICATION AND MEASUREMENT SETUP

This chapter describes the design, fabrication, and measurement methods used throughout this thesis. We outline the device design workflow based on electromagnetic simulations. The fabrication section presents the cleanroom processes used to realize superconducting devices, including thin-film deposition, electron-beam lithography, dry and wet etching, and double-angle shadow evaporation for Josephson junctions. Josephson junction fabrication is described for both Dolan-bridge and Manhattan junctions, along with the methods used for room-temperature junction characterization. Finally, we describe the device packaging and the cryogenic measurement setup. Together, this chapter aims to provide the practical details and cleanroom process considerations that are required to reproducibly fabricate and measure the devices presented in the following chapters.

2.1. DEVICE DESIGN

The superconducting devices presented in this thesis are often the outcome of an iterative design process that includes circuit design, electromagnetic simulation, nanofabrication, and cryogenic characterization. Small changes in device dimensions, material properties and fabrication-related imperfections can have a large impact on circuit parameters, so device design is often not a single-shot process. Instead, multiple design and fabrication iterations may be required to converge to desired resonance frequencies, coupling strengths, and loss characteristics.

The design process typically begins by specifying target circuit parameters, such as resonance frequencies, external coupling rates, and anharmonicities. These specifications are then mapped onto circuit elements, including capacitances and inductances, which in turn determine the physical layout of the device. As an initial guide, circuit quantization and modeling are useful tools for relating the target parameters to effective capacitances and inductances. Software such as QuCAT can model superconducting circuits in a lumped-element description and estimate quantities such as mode frequencies and anharmonicities [1]. Early in my PhD, device layouts were generated using stCAD [2], a custom Python-to-GDS design tool developed in our group, and simulated with SONNET. Later work made use of Quantum Metal [3] (formerly Qiskit Metal) interfaced with ANSYS electromagnetic solvers, providing a rapid feedback loop between layout design and electromagnetic simulation.

In a typical design workflow, the physical device layout is parametrically defined in Quantum Metal. The layout is then simulated using ANSYS Q3D or HFSS to extract electromagnetic properties such as resonance frequencies, coupling strengths, capacitances, and field distributions [4]. To ensure that the extracted parameters are reliable, the simulation setup is iterated until key quantities are converged. The resulting electromagnetic solution can be interpreted using either the lumped oscillator model (LOM) [5] or the energy participation ratio (EPR) method to connect the simulated fields to an effective circuit description [6]. LOM is typically used for rapid, order-of-magnitude design iterations and for extracting effective lumped parameters such as capacitances and inductances. The EPR method is used when more accurate estimates of nonlinear circuit parameters, such as anharmonicities and dispersive shifts, are required. Based on these extracted parameters, the circuit geometry is updated and the design is re-simulated until the target specifications are met. Once a satisfactory design is obtained, the layout is exported as a GDS file for fabrication.

For superconducting materials with a significant kinetic inductance contribution, such as NbTiN, kinetic inductance must be included in the design and simulation process to obtain accurate device parameters. When required, this is implemented using a modified fork of Quantum Metal that includes kinetic inductance modeling capabilities [7]. For NbTiN devices, neglecting kinetic inductance can for example lead to resonance frequencies that differ from measured values by approximately 20%. In contrast, for devices fabricated from aluminum or tantalum, the kinetic inductance contribution is much smaller and is therefore neglected in the design stage. It should be noted that Josephson junctions are not included in these simulations. Instead, during design and simulation they are represented as linear inductive elements within the circuit model. The physical junction layout is added at the final stage of the design process, where a predefined junction layout is inserted into the GDS file used for fabrication.

2.2. FABRICATION

This section describes the fabrication techniques used to produce the superconducting devices presented throughout this thesis. Since the performance of superconducting quantum circuits is highly sensitive to material quality, interfaces, and processing conditions, a clear and detailed description of the fabrication flow is essential to ensure reproducibility and to enable meaningful comparison between devices. The focus of this chapter is on the core fabrication steps that are common to all devices studied in this work, including substrate preparation, thin-film deposition, electron-beam lithography, etching, and shadow evaporation for Josephson junctions. Rather than presenting device-specific recipes in full detail, we establish a consistent fabrication framework that is referenced and adapted in subsequent chapters when needed. All devices presented in this work were fabricated at the Kavli Nanolab Delft [8].

2.2.1. FABRICATION OF LINEAR DEVICES

SUBSTRATES

The first step of the fabrication process is the selection of a suitable substrate material. Most of the devices presented in this thesis are fabricated on ultrapure, intrinsic $\langle 100 \rangle$ -oriented silicon wafers with a high resistivity ($\rho > 2.0 \times 10^4 \, \Omega \cdot \text{cm}$). In some cases, sapphire (Al_2O_3) substrates are used. Sapphire is a transparent, highly insulating material that can create some extra fabrication challenges, such as charging effects during lithography. However, its mechanical hardness, chemical inertness, and lower dielectric loss compared to silicon can make it useful for minimizing losses in superconducting quantum devices.

SUPERCONDUCTING THIN FILMS

The properties of superconducting thin films are strongly influenced by both the substrate and the deposition technique, including specific process parameters. The substrate not only provides mechanical support but also plays a critical role in the nucleation and crystal orientation of the superconducting film. Key factors such as surface roughness, chemical compatibility, and lattice mismatch between the substrate and the superconducting film can significantly impact the film quality. Common superconducting materials used in superconducting quantum circuits include Nb, Al, TiN, NbTiN and Ta [9–11]. In this work, superconducting films were deposited using two different physical vapor deposition (PVD) techniques.

The device described in chapter 3 used a 200 nm thick NbTiN film deposited on a 4-inch double side polished (DSP) silicon wafer via sputter deposition. Prior to deposition, the silicon wafer was cleaned using the RCA standard cleaning procedure. It was then transferred into the LLS801 DC magnetron sputtering system, a refurbished and customized industry-grade Evatec LLS801 system operated by the Space Research Organization Netherlands (SRON) [12]. During deposition, a flow of nitrogen gas was introduced into the chamber, reacting with a 99.95 % pure NbTi target to form a stoichiometric NbTiN film on the substrate.

The device described in chapter 5 used a 100 nm thick Tantalum film deposited on a 4-inch double side polished (DSP) Sapphire wafer. Tantalum was used for its low microwave loss in the α -phase, which has been associated with significant improvements in qubit coherence times [13]. The tantalum-coated wafers used in this work were sourced from STARCryo.

WAFER DICING AND CHIP PREPARATION

After thin-film deposition, some of the wafers were divided into individual chips for further processing. For most devices presented in this thesis, including resonators and qubits, dicing was performed directly after the superconducting film was deposited. In contrast, the SNAIL-based parametric amplifiers described in Chapter 4 were fabricated entirely at the wafer scale, and were only diced after all fabrication steps were completed.

Wafer dicing was performed using a Disco DAD 3220 automatic dicing system, shown in Fig. 2.1 a. The final chip dimensions depend on the device type and PCB requirements. The wafers for resonator and qubit devices were diced into square chips of $10 \times 10 \text{ mm}^2$, while wafers containing SNAIL-based parametric amplifiers were diced into rectangular chips of $11 \times 5 \text{ mm}^2$, as shown in Fig. 2.1 b. Prior to dicing, a protective resist layer (either S1805 or PMMA) was applied to the wafer surface to reduce particle contamination and mechanical damage during cutting. For wafers that already contained Josephson junctions, special care was taken during this step, as elevated baking temperatures can lead to small changes in junction parameters. In these cases, the resist was baked at reduced temperatures or in a vacuum oven.

Silicon wafers were diced using resin-bonded diamond blades (typically NBC-ZB 1050 or Z09-SD2000) at a rotation speed of 30 000 RPM. Sapphire wafers required modified dicing conditions due to their higher mechanical hardness. For these wafers, a different blade (VT07-SD320-VC100-75) was used together with a reduced rotation speed of approximately 15 000 RPM. The cut depth was also increased in incremental steps, starting with a shallow cut and followed by successive deeper cuts with adjusted feed speeds. After dicing, individual chips were cleaned to remove residual resist and particulates before further processing or packaging.

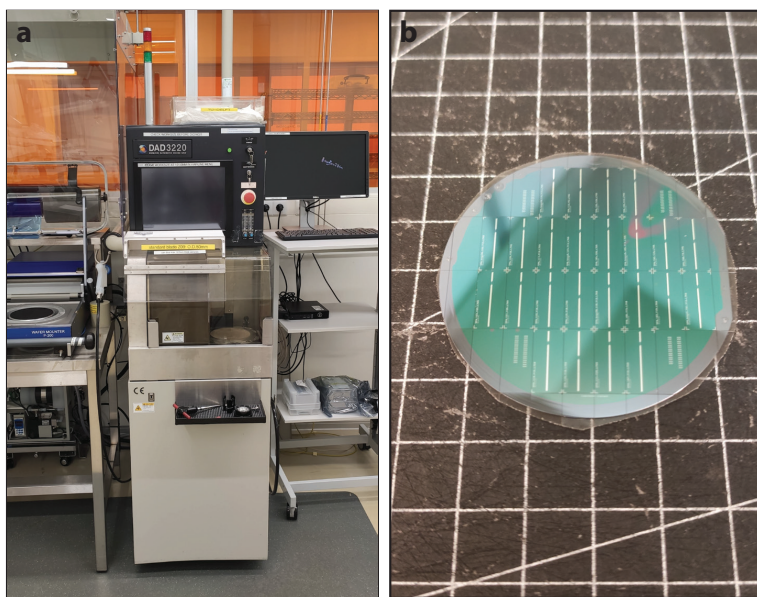


Figure 2.1: Wafer dicing into individual chips. (a) Photograph of the Disco DAD 3220 dicing system. **(b)** A 2-inch wafer after dicing into individual $11 \times 5 \text{ mm}^2$ chips.

LITHOGRAPHY

Once a superconducting film has been deposited, the next step in the fabrication process is to define the device geometry. Mask-based optical lithography is widely used in industrial microfabrication due to its high throughput and reproducibility. However, for academic research, where rapid design iterations and low-volume fabrication are often required, direct-write lithography techniques are more suitable. In the Kavli Nanolab Delft, both a direct-write optical laser writer and electron-beam lithography (EBL) systems are available. For the devices presented in this thesis, electron-beam lithography was used exclusively as it provides the highest resolution and reproducibility for superconducting quantum devices. An overview of the electron-beam lithography system is shown in Fig. 2.2.

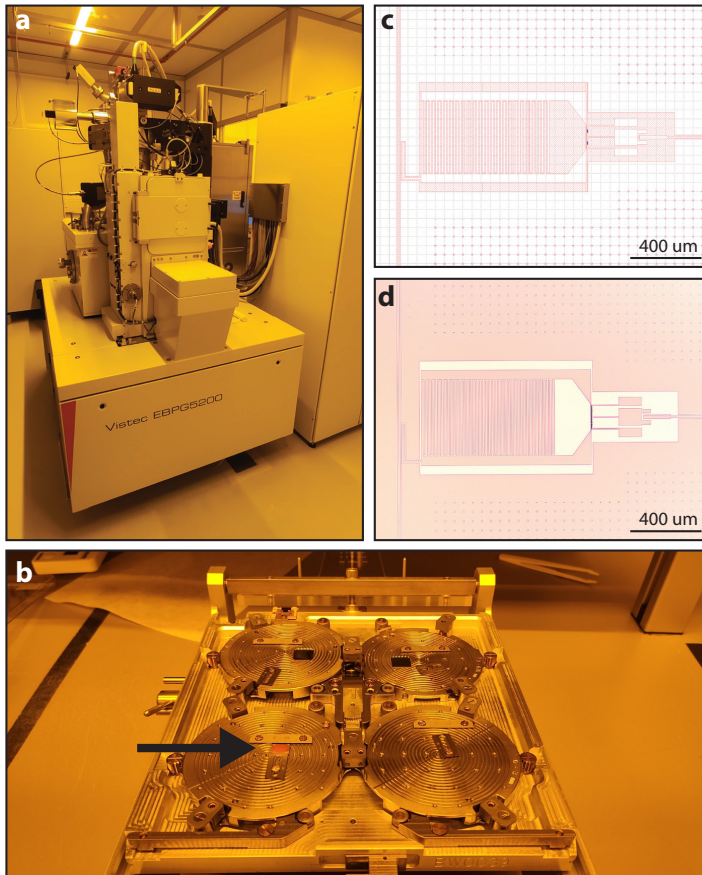


Figure 2.2: Electron-beam lithography system. (a) Raith EBPG 5200 electron-beam lithography system used for device patterning at the Kavli Nanolab Delft. (b) A $10 \times 10 \text{ mm}^2$ chip is mounted on the EBPG sample holder using clamps to ensure electrical grounding during exposure. (c) Example GDS layout file defining the lithographic pattern of a superconducting device. (d) Optical microscope image of the same design on a chip after electron-beam exposure and resist development, demonstrating the pattern transfer from the GDS layout to the chip.

ELECTRON-BEAM LITHOGRAPHY (EBL)

Lithography was performed using a Raith EBP 5200 or EBP 5200+ electron-beam lithography system. Fig. 2.3 shows a schematic process flow for transferring a pattern to the substrate via etching. Before ebeam exposure, the substrate is spin-coated with an electron-beam resist, such as Polymethyl Methacrylate (PMMA) or AR-P 6200.18 (CSAR 62), to form a uniform resist layer. Of which the final thickness depends on the spin coating rotation speed and the resist properties. The required resist thickness depends mainly on later processing steps. To obtain the highest resolution, it is best to work with thin resist layers. However, a thicker resist layer might be necessary when it is used as an etching mask. Following spin coating, the resist-coated substrate is then soft-baked on a hot plate, which removes solvents to solidify the liquid resist film.

During the EBL exposure, a focused electron beam is used to directly write patterns onto the resist, which effectively modifies the solubility of the resist in the exposed areas. After exposure, an optional post-exposure bake can be performed, to harden the resist in the case of follow up etching steps. The samples can then be developed with an appropriate developer for the specific resist. This removes the exposed resist in the case of positive tone resist, or the unexposed area for negative tone resist. A resist mask is then formed that can be used for further processing steps, such as reactive ion etching (RIE) or metal deposition and lift-off.

We note that special care should be taken when working with insulating substrates such as sapphire. During electron-beam exposure, substrate charging can lead to beam deflection, pattern distortion, and even potential damage the EBP. To mitigate these effects, it is essential that the top of the sample surface is electrically conductive and connected to the EBP holder during exposure. Even when a metallic thin film is already present on an insulating substrate, a thin conductive polymer layer such as AR-PC 5090 (Electra 92) is applied on top of the resist stack to dissipate electrical charges during exposure and ensure stable exposure conditions.

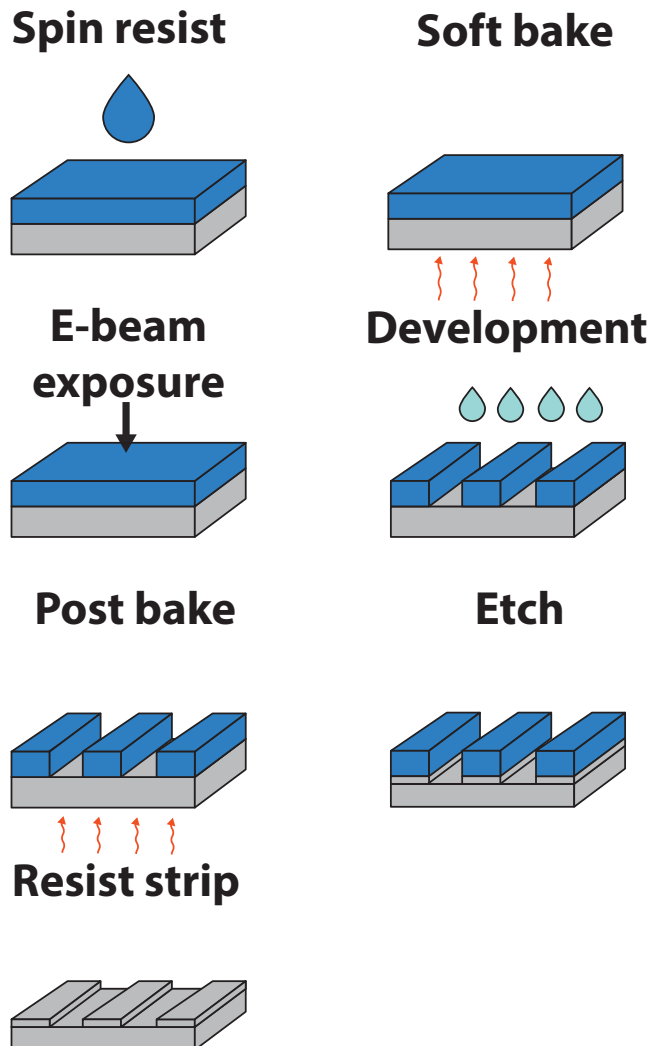


Figure 2.3: Schematic electron-beam lithography (EBL) process flow for pattern transfer by etching. The substrate is spin-coated with an electron-beam resist, followed by a soft bake on a hot plate to remove the solvents. The pattern is written by e-beam exposure, and the resist is developed afterward to form the lithographic mask. An optional post-exposure bake can be used to harden the resist prior to dry etching. The pattern is then transferred into the underlying substrate by etching. Finally, the resist is stripped.

ETCHING

Etching is used to (selectively) remove material or transfer a lithographically defined pattern onto the underlying substrate or thin film. The choice of etching technique is determined by the required etch profile, feature size, surface roughness, and the sensitivity of the material to different etching processes. A key distinction in etching processes is between isotropic and anisotropic etching. In isotropic etching, material is removed equally in all directions, resulting in rounded profiles and an undercut will be formed under the resist mask. This can limit the use in nanoscale device fabrication when the size of the undercut becomes comparable to the desired feature size.

In contrast, anisotropic etching is highly directional, with material being removed primarily perpendicular to the substrate surface. This enables the transfer of the exact mask dimensions into the substrate. Vertical sidewalls are formed in the etching profile that directly represent the lithographic pattern, which is important to obtain well-defined nanoscale devices. An illustration of these two etching regimes is shown in Fig. 2.4.

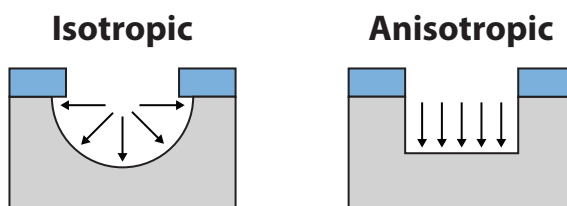


Figure 2.4: Schematic illustration of ideal isotropic and anisotropic etching. In isotropic etching, the etch rate is the same in all directions, producing an undercut beneath the resist mask. In anisotropic etching, the etch is highly directional, yielding vertical sidewalls.

Etching techniques are broadly classified into two broad categories: wet etching and dry etching. In this work, dry etching was mainly used for the definition of device features, while wet etching was used for substrate cleaning and oxide removal, as well as in some post dry etching steps to improve surface quality.

DRY ETCHING

Dry etching is a material removal process without the use of liquid chemicals. Several techniques fall into this category, including plasma etching, ion beam milling, and reactive gas etching. These methods remove material through chemical reactions with reactive gasses, physical etching by energetic particles, or through a combination of both.

Device patterning was performed using plasma-based Reactive Ion Etching (RIE). The working principle of RIE is illustrated in Fig. 2.5. In this process, the sample is placed on one of two parallel electrodes inside a vacuum reaction chamber. After the chamber is pumped down to a low base pressure, a controlled flow of process gas such as sulfur hexafluoride (SF_6) or Carbon tetrafluoride (CF_4) is introduced into the chamber. Applying radio-frequency (RF) power between the electrodes ionizes the gas and creates a plasma containing chemically reactive radicals and energetic ions.

The reactive radicals diffuse to the substrate surface, where they react with the exposed ma-

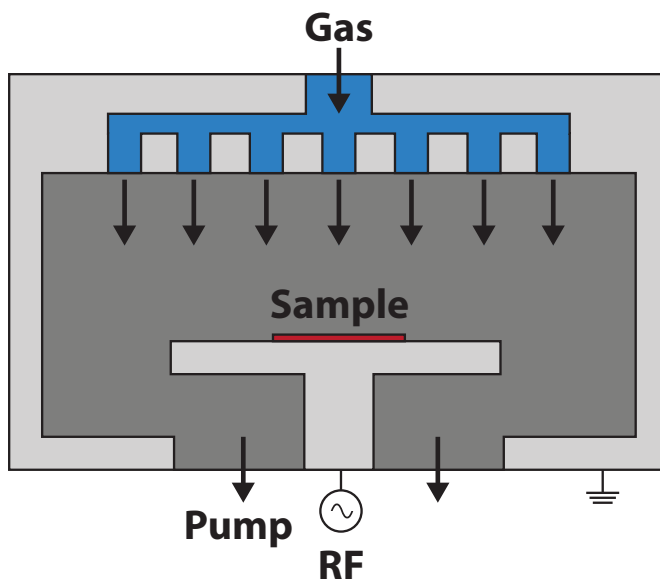


Figure 2.5: Schematic illustration of the working principle of reactive ion etching (RIE). The sample is placed on an electrode inside a vacuum chamber. Process gas is introduced and RF power is applied between the parallel electrodes to generate a plasma. Reactive radicals reach the surface and form volatile reaction products, which are removed by the vacuum pump. Simultaneously, ions are accelerated toward the sample, making the etch more anisotropic.

terial to break chemical bonds and form volatile reaction products. These by-products desorb from the surface and are continuously removed by the vacuum pump. Simultaneously, ions are accelerated toward the sample surface by the electric field and contribute to physical sputtering. By adjusting process parameters such as gas composition, chamber pressure, RF power, and bias voltage, the relative contributions of chemical and physical etching can be tuned, allowing control over the resulting etch profile.

For metal etching, endpoint detection is often used to determine when the target film has been fully removed. In this work, endpoint detection is performed optically by monitoring changes in the reflected signal from the wafer surface using a laser. Because the optical reflectivity of the metallic film differs from that of the underlying substrate, a clear change in the reflected signal indicates when the metal layer has been removed. Once the endpoint is reached, the etch is typically continued for an additional 10 seconds to ensure complete removal of the film across the entire sample. Longer over-etch times are avoided whenever possible, as both silicon and, to a much lesser extent, sapphire substrates can be etched or damaged under these conditions.

During the etching process, the photoresist mask is also gradually etched by the plasma. To ensure accurate pattern transfer, a sufficiently thick resist layer must therefore be used to withstand the full etch duration. After completion of the etch, a short in-situ oxygen plasma clean is performed before breaking vacuum to remove fluorocarbon residues that may form on the sample or chamber surfaces. Without this oxygen plasma step, an insulating residue was

occasionally observed on top of NbTiN films, which can negatively affect subsequent processing steps. While reactive ion etching provides precise control over the feature dimensions and sidewall profiles, it is not always optimal for obtaining a smooth surface profile or when highly selective material removal is desired. In these situations, wet etching can be used either as an alternative to dry etching or as a complementary post-processing step, as discussed in the next section.

WET ETCHING

Wet etching removes material through chemical reactions in a liquid etchant. The etching rate is mainly determined by the concentration of the etchant, the bath temperature, and the transport of reactants to and from the surface. Wet etching processes can be divided into anisotropic and isotropic etching.

Anisotropic wet etching is achieved using selective chemical solutions that remove material at different rates depending on crystallographic orientation. A well-known example is single-crystal silicon, which can be etched in potassium hydroxide (KOH) or tetramethylammonium hydroxide (TMAH), where the $\langle 100 \rangle$ crystal planes etch faster than other planes. This property is widely exploited in the fabrication of micro-electromechanical systems (MEMS) [14].

In contrast, isotropic wet etching removes material at approximately the same rate in all directions, independent of the crystallographic orientation. While isotropic wet etching can be used to define lithographic patterns, the inherent undercut can limit the use for high-resolution or high-aspect-ratio nanoscale features. In this work, isotropic wet etching was therefore primarily used for surface cleaning and surface preparation following lithography, as well as for short etch steps after dry etching to improve surface quality. The wet-etching processes most frequently used in this work include:

- **RCA standard clean:** A multi-step wet chemical cleaning process originally developed by the Radio Corporation of America (RCA), used to remove organic residues, metal ions, and native oxides from silicon wafers [15]. The RCA standard clean was applied prior to the deposition of NbTiN films. In addition, the first step of the RCA clean (SC-1) was occasionally used as a final cleaning step after dry etching NbTiN to obtain a smooth surface with minimal contamination.
- **Piranha solution ($\text{H}_2\text{SO}_4/\text{H}_2\text{O}_2$):** A strong oxidizing solution used to remove organic residues from the sample surface.
- **Hydrofluoric acid (HF):** Used for oxide removal, for example to strip native silicon oxide (SiO_2) and to remove surface oxides on metal films, while etching the underlying silicon much more slowly. In practice, a buffered oxide etch (BOE) was often used instead of pure HF to obtain a slower, more controllable etch rate. Care must be taken to ensure compatibility with previously deposited metals, as commonly used superconductors such as aluminum are etched by HF.

By using a combination of the fabrication techniques described above, it is possible to fabricate linear superconducting devices, such as the coplanar-waveguide resonators presented in Chapter 5.

2.2.2. FABRICATION OF NON-LINEAR DEVICES

Building upon the fabrication processes developed for linear devices, we now describe the fabrication of nonlinear circuit elements based on Josephson junctions. These junctions provide the nonlinearity required for superconducting qubits and parametric amplifiers. In this thesis, superconductor–insulator–superconductor (SIS) Josephson junctions of the Al–AlO_x–Al type are used, fabricated by double-angle shadow evaporation.

Two junction architectures are used: Manhattan-style junctions for qubits and Dolan-bridge junctions for SNAIL-based parametric amplifiers. While the overall fabrication flow is similar for both approaches, specific parameters such as resist thickness, electron-beam exposure dose, and evaporation angles differ. In the following, we first describe the common fabrication steps, followed by details specific to each junction architecture.

DOUBLE LAYER RESIST STACK

To define the Josephson junction area and to enable a reliable lift-off process, we use a bi-layer MMA/PMMA resist stack. The bottom layer consists of MMA(8.5) MAA copolymer, while the top layer is PMMA 950k. A schematic of the developed resist profile is shown in Fig. 2.6 (a). The MMA copolymer layer has a higher sensitivity to electron-beam exposure than the PMMA layer. After development, this difference in sensitivity results in an undercut resist profile, which is essential for a clean lift-off after metal evaporation. For some junction designs, an additional low-dose “ghost” exposure was applied in a narrow region surrounding the junction area. This creates an enlarged undercut, as illustrated in Fig. 2.6 (b), which can improve the lift-off process and reduce the risk of metal deposited on the resist sidewalls forming shorts or residual fences during angled evaporation.

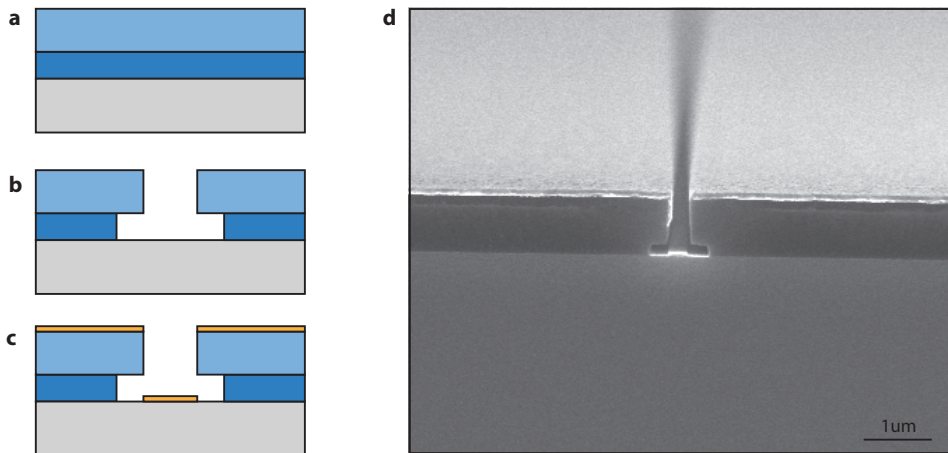


Figure 2.6: Double layer resist stack. (a–c) Schematic of the resist stack at different stages: (a) before development, (b) after electron-beam exposure and development, and (c) after metal evaporation. The undercut ensures that a clean lift-off process can be achieved. (d) Cross-section SEM of the resist stack after angled metal evaporation. The large undercut was obtained by an additional low dose exposure around the junction area.

DOUBLE-ANGLE EVAPORATION AND LIFT-OFF

Josephson junctions were fabricated by double-angle evaporation in a Plassys MEB 550 electron-beam evaporation system. Figure 2.7 shows the evaporation system used in this work, along with a close-up view of a chip mounted on the angled sample holder during deposition. The working principle of the evaporation system is illustrated schematically in Fig. 2.8. After loading the sample into the evaporation chamber, the first aluminum layer was deposited at a predefined angle relative to the substrate normal, forming the bottom electrode of the junction.

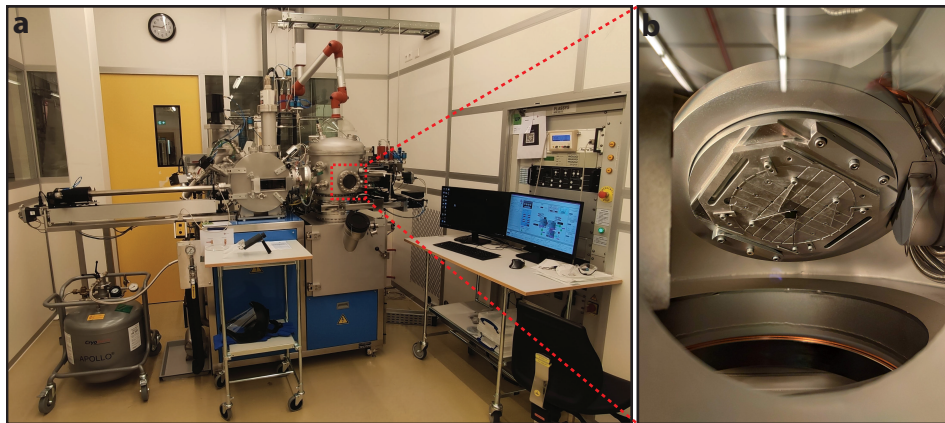


Figure 2.7: The Plassys electron-beam evaporation system used for Josephson junction fabrication. (a) Photograph of the Plassys MEB 550 electron-beam evaporation system installed in the cleanroom. The system features a rotatable, tilt-adjustable sample stage and supports in-situ oxidation. **(b)** Close-up view of a 10×10 mm chip mounted on the sample holder inside the main evaporation chamber during deposition. The sample is tilted relative to the substrate normal to enable angled shadow evaporation.

Following the first deposition, oxygen was introduced into the chamber to oxidize the aluminum surface and form the AlO_x tunnel barrier. The oxidation pressure and duration were chosen to set the desired junction resistance. After oxidation, a second aluminum layer was deposited at a different angle to form the counter-electrode. The overlap between the two aluminum layers defines the effective Josephson junction area. Depending on the junction architecture, this second deposition is performed either from the opposite direction for Dolan-bridge junctions or after a 90° rotation of the sample for Manhattan junctions. These approaches are discussed separately in the following sections.

When a Josephson junction requires a galvanic connection to a previously defined superconducting film, such as NbTiN or tantalum, the native oxide on the exposed metal must be removed prior to aluminum deposition. Two different approaches were used to achieve this. In the first approach, the sample was briefly dipped in a buffered oxide etchant (BOE) immediately before loading into the Plassys evaporation system. This removes the native oxide from the exposed superconducting film, after which the sample is rapidly transferred into vacuum to minimize reoxidation. Alternatively, in-situ argon ion milling was performed in the load-lock of the Plassys to remove the oxide layer. While this method is effective in removing the oxide, argon milling can also damage the exposed substrate surface, which may cause ad-

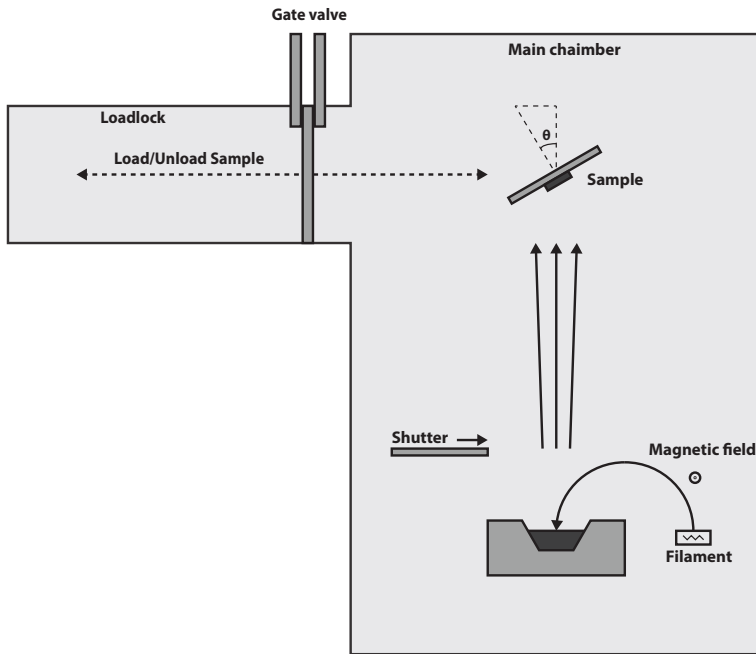


Figure 2.8: Schematic illustration of the Plassys MEB 550 electron-beam evaporation system. The system consists of a load-lock chamber for sample loading, unloading, and optional in-situ argon ion milling, and a main deposition chamber. The sample is mounted on a tiltable and rotatable holder. After transfer into the main chamber, this allows for deposition at a controlled angle θ relative to the substrate normal. An electron beam locally heats the source material to generate a vapor flux, while a mechanical shutter defines the start of deposition and enables precise thickness control.

ditional loss. To mitigate this, the milling could be performed at a large angle relative to the substrate normal while rotating the sample, such that the junction area itself remained mostly protected by the resist stack.

Following metal deposition, the unwanted metal and remaining resist were removed using a lift-off process. Samples were typically immersed in acetone heated to 50 °C for several hours to dissolve the resist. A pipette was then used to gently remove the detached metal film. Brief sonication was also often applied afterwards to assist lift-off. Finally, the samples were rinsed in isopropanol and dried with a nitrogen gun.

DOLAN BRIDGE JUNCTIONS

The Dolan-bridge technique, first introduced by Dolan in 1977 [16], is one of the two widely used methods for fabricating Josephson junctions by shadow evaporation. In this approach, a narrow resist bridge is defined by electron-beam lithography and suspended above the substrate, allowing two angled metal evaporations to overlap only in a well-defined region.

Figure 2.9 illustrates the principle of Dolan-bridge shadow evaporation. The schematic shows both the fabrication of a single Josephson junction as well as multiple junctions connected

in series using multiple suspended bridges. After lithography and development, the first aluminum layer is deposited at an angle, forming the bottom electrode. This layer is then oxidized in situ to create the tunnel barrier. A second aluminum layer is subsequently deposited from the opposite angle, forming the top electrode. After lift-off, the overlap between the two metal layers defines the Josephson junction area.

When fabricating multiple Josephson junctions in series using the Dolan-bridge technique, the junction areas are set by different shadowing mechanisms for the outer and central junctions. For the outer junctions, the overlap of the two angled evaporations is defined by shadowing from the underside of each suspended resist bridge. In this case, the junction area depends on the bridge width, resist thickness, and evaporation angle, which together determine the extent of the shadow cast by the bridge during angled deposition.

In contrast, the central junction is defined by shadowing from the upper edges of two neighboring resist bridges. Here, the overlap region is determined by the spacing between the adjacent bridges, as the two angled evaporations overlap on the substrate between them, without relying on shadowing from a suspended underside. As a result, the central junction area is primarily set by the bridge separation and is less sensitive to the exact suspension quality of the bridge.

Achieving uniform junction areas across all junctions therefore requires separate control over these two mechanisms. In practice, the electron-beam exposure and development conditions are first optimized to produce well-suspended bridges with the desired geometry, ensuring reproducible outer junctions. The spacing between neighboring bridges is then adjusted to tune the central junction area such that it matches the outer junctions.

Because the Dolan-bridge technique relies critically on the formation of a fully suspended resist bridge with a well-controlled undercut, even small deviations in the lithography dose, development conditions, or resist thickness can lead to fabrication failures. Imperfect suspension, insufficient undercut, or mechanical instability of the bridge can compromise the shadowing mechanism required for reliable junction definition.

These fabrication issues usually appear in characteristic repeatable patterns that can be identified in scanning electron microscope (SEM) images. As such, SEM inspection provides a powerful diagnostic tool for assessing process failures and for guiding the optimization of lithography and development parameters.

Figures 2.10 and 2.11 present examples of both successful and failed Dolan-bridge junctions observed during fabrication. For each case, SEM images are shown alongside schematic cross sections that illustrate the underlying resist profile responsible for the observed outcome. Together, these examples illustrate common failures in Dolan-bridge shadow evaporation and clarify how Dolan-bridge imperfections result in junction defects.

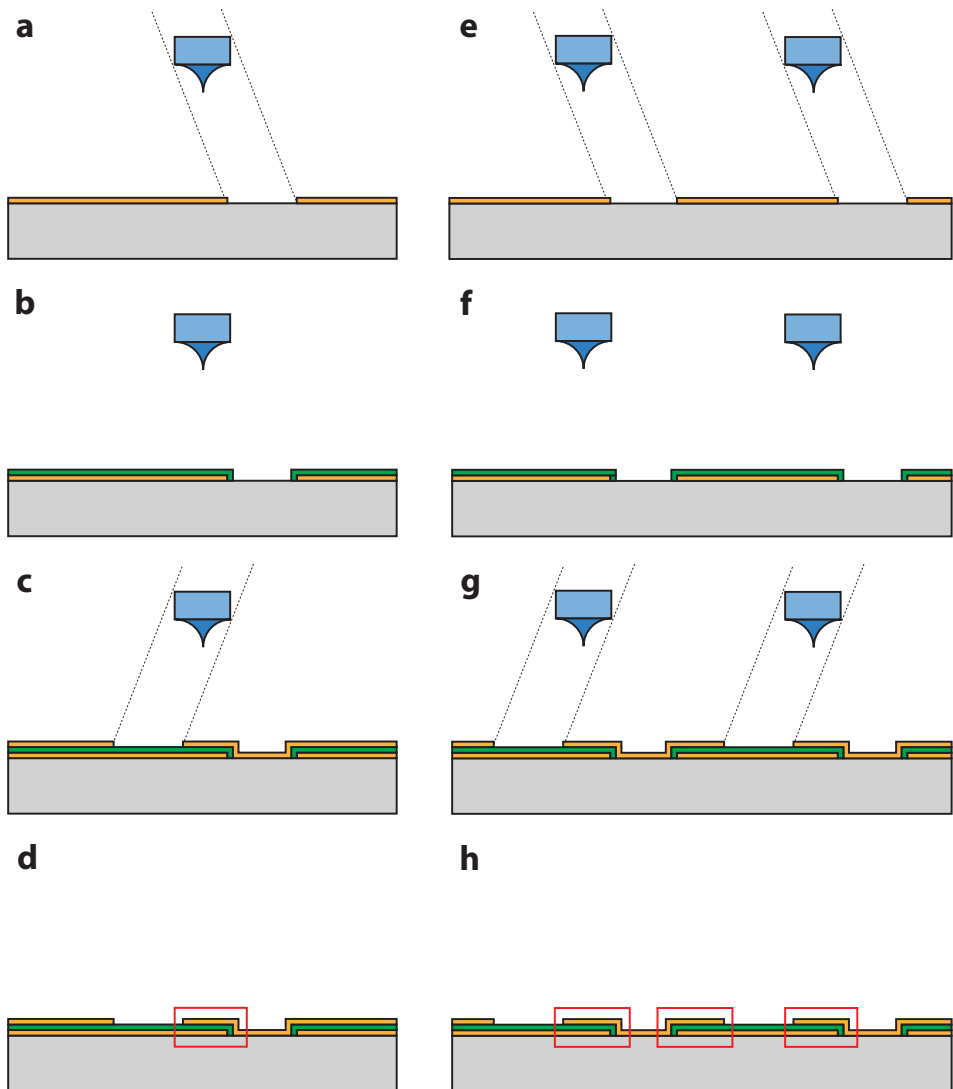


Figure 2.9: Schematic illustration of shadow evaporation using the Dolan Bridge technique. (a-d) shows the fabrication of a single Josephson junction, while (e-h) shows the fabrication of three junctions connected in series. (a,e) After electron-beam lithography and development, the first metal layer is deposited at an angle, forming the bottom electrode. (b,f) The deposited metal is oxidized to form the tunnel barrier, indicated in green. (c,g) A second metal layer is deposited from the opposite angle, forming the top electrode with a controlled overlap area. (d,h) After lift-off, the overlap regions highlighted by red boxes define the final Josephson junctions.

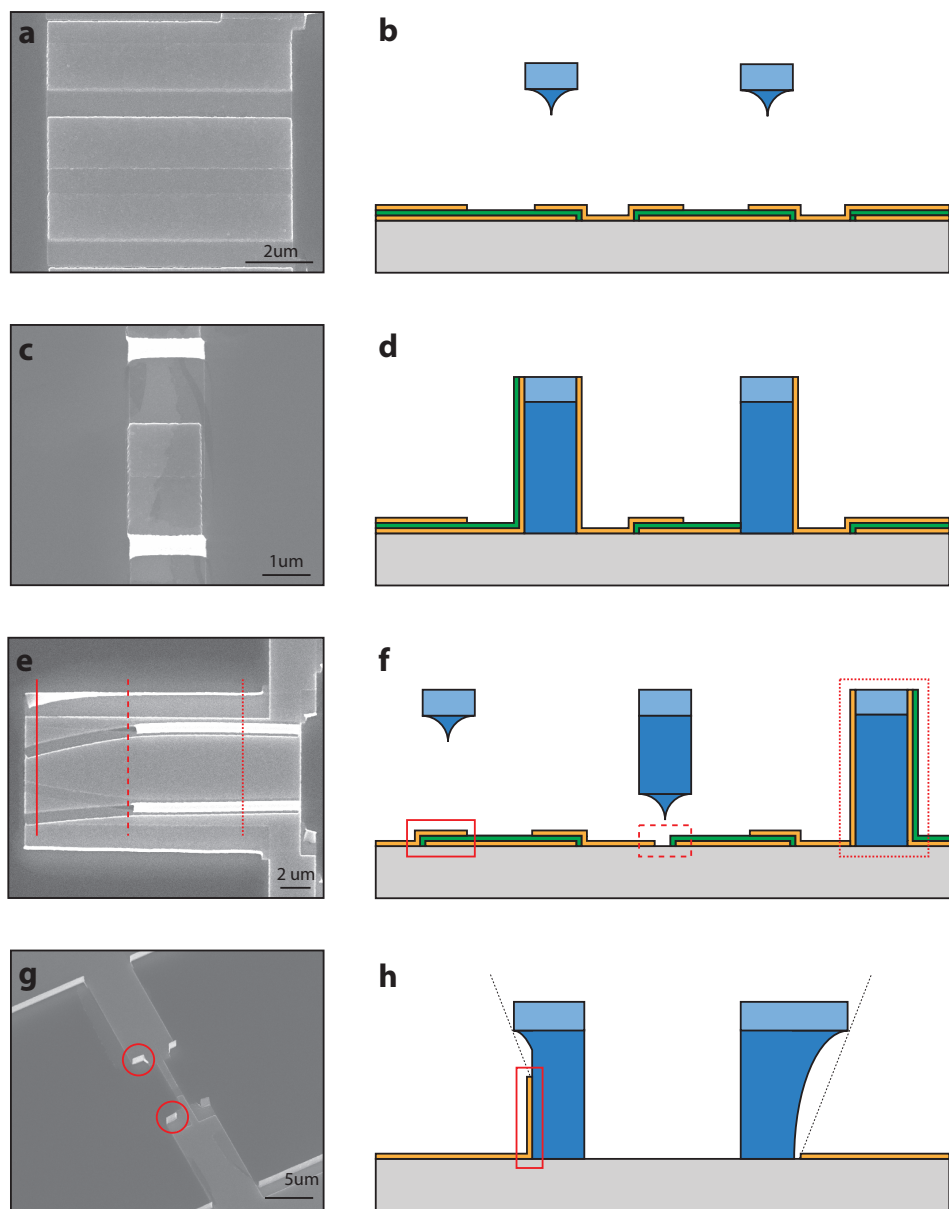


Figure 2.10: Common success and failures in Dolan-bridge shadow evaporation. (a,b) Successful Josephson junction formation using a fully suspended resist bridge with sufficient undercut. (c,d) A resist bridge that is not suspended. During shadow evaporation, metal is deposited on the resist sidewalls, preventing proper junction formation at the outer two locations. The central junction remains well defined because it is formed by the top of the resist profiles and does not rely on shadowing from the underside of a suspended bridge. (e,f) A partially broken or partially suspended resist bridge, leading to incomplete or asymmetric junction formation. (g,h) Insufficient undercut at the ends of a large exposed region allows metal to deposit onto the resist sidewalls during evaporation. This can be avoided by ensuring that there is a sufficiently large resist undercut.

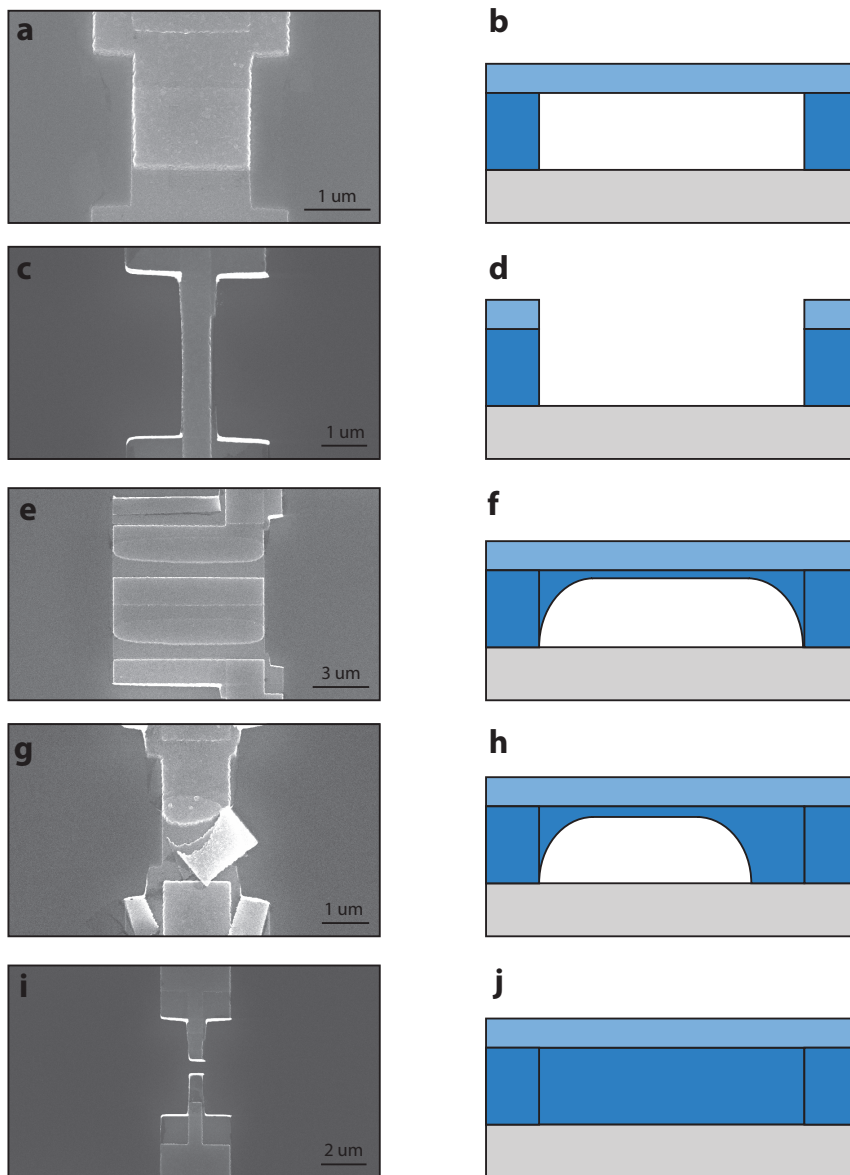


Figure 2.11: Additional success and failures in Dolan-bridge lithography. (a,b) A well-defined and fully suspended resist bridge with sufficient undercut, suitable for reliable shadow evaporation. (c,d) A missing resist bridge, either due to overexposure or mechanical failure during development. In this case, junctions will not be formed. (e,f) Insufficient exposure at the ends of the bridge results in rounding of the resist profile. This reduces the effective undercut and compromises shadowing during angled evaporation. (g,h) Similar to panels (e,f), but asymmetric: one side of the bridge remains fully supported while the opposite side exhibits rounding. (i,j) The bridge is not suspended at all, preventing the formation of a tunnel junction and resulting in an open circuit connection.

MANHATTAN JUNCTIONS

For qubit fabrication, we use Manhattan-style Josephson junctions instead of Dolan-bridge junctions. This junction architecture was developed to improve fabrication robustness and reproducibility [17] and has since been widely adopted in the fabrication of superconducting qubits [18]. Manhattan junctions are formed by the overlap of two orthogonal electrodes deposited in successive angled evaporations with a 90 degree in-plane rotation between depositions, rather than by using a suspended resist bridge. As a result, the junction overlap area is less sensitive to resist variations. In particular, it is not affected by small changes in resist thickness. Furthermore, it does not rely on fragile suspended bridges, which often limits Dolan-bridge junctions [19]. This can improve device-to-device uniformity, which is particularly important for qubits, where small junction variations translate into shifts of the qubit frequency and anharmonicity [20, 21].

Figure 2.12 illustrates the fabrication sequence for Manhattan-style junctions. After patterning the MMA/PMMA resist stack, the first aluminum layer is deposited at an angle θ_1 , forming the bottom electrode. This layer is then oxidized in situ to create the tunnel barrier. The sample is then rotated by 90°, and a second aluminum layer is deposited at an angle θ_2 , forming the top electrode. The Josephson junction is then defined by the overlap of the two deposited metal layers.

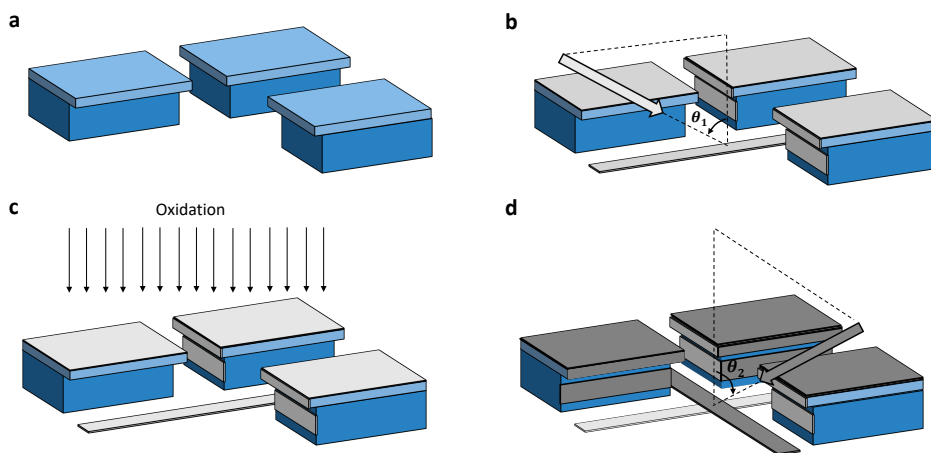


Figure 2.12: Schematic illustration of Manhattan-style Josephson junction fabrication by angled shadow evaporation. (a) Resist pattern defining the junction electrodes prior to metal deposition. **(b)** First metal evaporation at an angle θ_1 , forming the bottom electrode. **(c)** In-situ oxidation of the first metal layer to create the tunnel barrier. **(d)** Second metal evaporation at an angle θ_2 , performed after a 90° rotation of the sample, forming the top electrode. The Josephson junction is defined by the overlap of the two deposited metal layers.

The smallest and largest junction sizes that can be fabricated in this way are determined by the combination of the resist profile and the evaporation conditions. The resist thickness and undercut must be sufficient so that the metal deposited on the resist sidewall does not connect to the substrate, which can compromise lift-off and lead to shorts.

To determine the maximum reliable resist opening for the present resist stack and evaporation conditions, we inspect cleaved samples by cross-sectional SEM after evaporation. When interpreting these cross sections, it is important to note that SEM imaging itself can modify the resist profile.

Figure 2.13 shows that repeated imaging of the same region at increasing magnification can lead to progressive rounding and deformation of the suspended MMA/PMMA resist stack due to accumulated electron-beam exposure. Rounded resist features observed at high magnification should therefore often be treated as an imaging-induced artifact rather than as an issue with the resist profile.

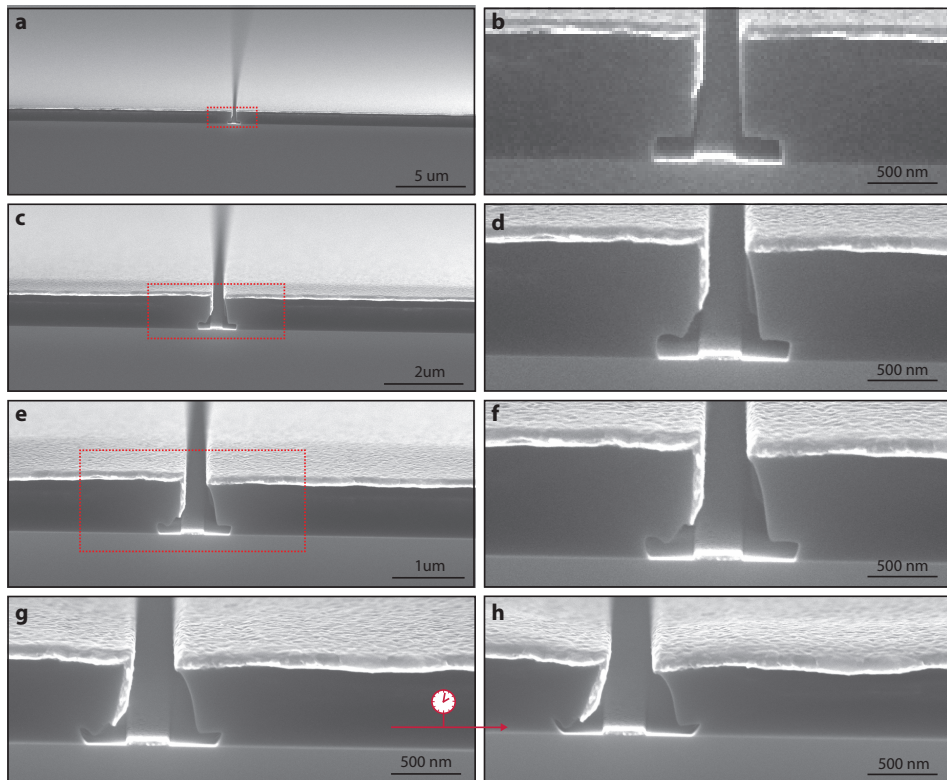


Figure 2.13: Electron-beam-induced deformation of a double-layer resist stack during SEM imaging. Cross-sectional SEM images of an MMA/PMMA resist stack after angled aluminum evaporation. Panels (a,c,e) show cross-sectional images acquired at increasing magnification (12 000 $\times$, 35 000 $\times$, and 65 000 $\times$), while panels (b,d,f) show cropped views of the regions indicated by red dashed boxes, displayed at a comparable scale. As the same region is imaged sequentially at increasing magnification, the accumulated electron-beam exposure increases, leading to progressive rounding and deformation of the suspended resist. Panels (g) and (h) are taken at the same magnification, with panel (h) recorded after additional consecutive SEM images, illustrating the continued resist deformation due to repeated electron-beam exposure.

With this consideration in mind, Fig. 2.14 compares cross-sectional SEM images of Manhattan junctions after angled aluminum evaporation for different junction widths. For a typical junction width used in this work, which is typically around 100 nm, the metal defines well-separated electrodes and lift-off remains clean. At 300 nm, near the upper end of reliable fabrication for the present resist stack, additional metal on the resist sidewalls is visible, but it does not connect to the substrate, such that lift-off is still successful. For very wide openings (800 nm), shadowing becomes insufficient and metal can reach the substrate in unintended regions, which may lead to electrical shorts after lift-off.

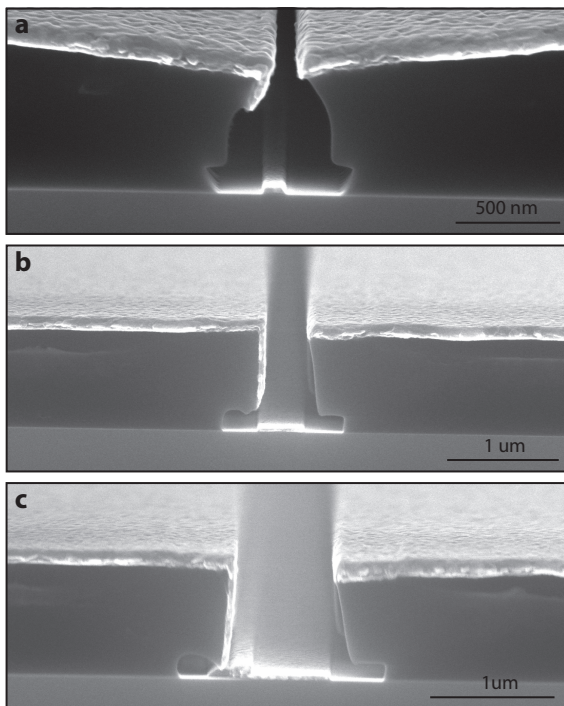


Figure 2.14: Cross-sectional SEM images of Manhattan junctions after metal evaporation for different junction widths. Panels (a–c) show junction cross sections with nominal widths of (a) 100 nm, (b) 300 nm, and (c) 800 nm, highlighting the effect of the resist opening width on the resulting metal profile and sidewall coverage.

Based on these observations, we typically target junction widths between 50 nm and 250 nm. The exact width limit is not universal and depends on the resist thickness and undercut, the evaporation angles, and the total deposited metal thickness. Increasing the resist thickness/undercut or adjusting the evaporation angle can extend the possible junction size range, whereas thicker metal films generally reduce it.

ROOM-TEMPERATURE JUNCTION PROBING

Probe junctions are included on each chip to characterize the normal-state resistance R_N across a range of junction sizes. Figure 2.16 shows an optical image of a typical probe junction array. These test structures enable rapid verification of junction quality and uniformity from room-temperature measurements. The array format allows efficient comparison of multiple junctions with different widths, providing a fast feedback loop during the fabrication process. The measurement is performed using a two-point probe configuration. Two tungsten probe needles are placed on the designated contact pads of a given test junction, and a small DC current bias of $I = 10$ nA is applied. The resulting voltage drop across the junction is measured,

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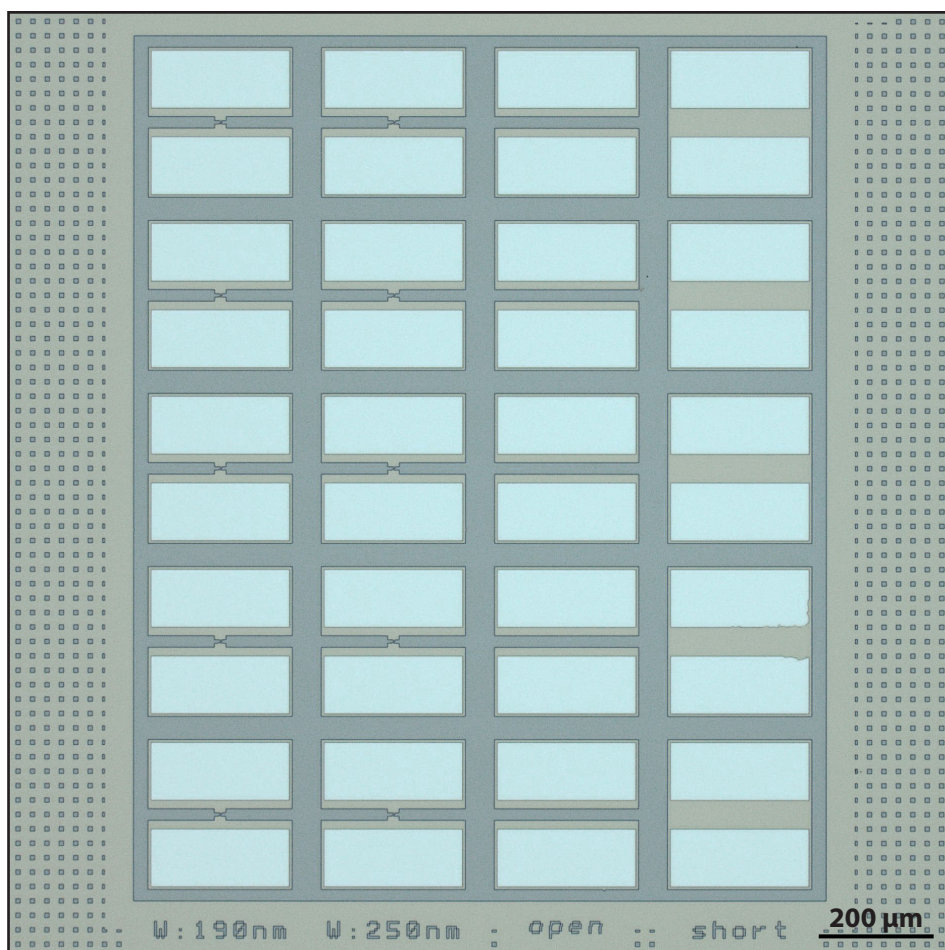


Figure 2.15: Arrays of probe junctions. Optical image showing an array of probe junctions used for measuring the normal state resistance. The Al junctions are located between two larger probe pads. For devices fabricated on NbTiN or Ta base layers, additional Al probing pads are placed on top of the pads to ensure reliable electrical contact with the probe needles.

from which the normal-state resistance $R_N = V/I$ can be obtained. Typical values of R_N lie in the k Ω range, depending on the junction area and the thickness of the oxide barrier.

To improve measurement accuracy, both open and shorted test pads are included in the layout. The open pads allow verification that the etch process has fully removed the base metal in regions intended to be isolated. The shorted pads, on the other hand, can be used to determine the contact resistance associated with the probe-pad interface. By subtracting this contact resistance, a more accurate estimate of the intrinsic junction resistance can be obtained. For the highest precision measurement, a four-probe measurement configuration is preferred. In this arrangement, separate pairs of contacts are used to source current and to measure the voltage drop, effectively eliminating both contact and lead resistances. When shorted junctions are then also patterned with the same geometry as the measured junctions, this also accounts for the series resistance of the leads between the pads and the junction area. As precise junction targeting was typically not our main priority, we primarily used fast and simple two-probe measurements.

It is important to note that the probe station's light source can influence the measured resistance, especially when using silicon substrates. Illumination may generate charge carriers in the substrate, thereby modifying the measured junction resistance. Consequently, the most accurate measurements are obtained with the light source turned off. Furthermore, one should take care not to switch the light on or off while the probe needles are touching the device. The sudden change in illumination can briefly drive unwanted currents that may destroy the junctions.

Estimating Josephson parameters from room-temperature measurements.

The normal-state resistance R_N obtained from these room-temperature tests provides a way to estimate several low-temperature junction parameters, such as the critical current I_c , the Josephson inductance L_J , and the Josephson energy E_J . These quantities are related through the Ambegaokar-Baratoff relation [22], which describes the relation between the critical current I_c and the normal-state resistance R_N for an ideal Josephson junction formed between two identical superconductors:

$$I_c R_N = \frac{\pi \Delta(T)}{2e} \tanh\left(\frac{\Delta(T)}{2k_B T}\right) \quad (2.1)$$

where e is the elementary charge, $\Delta(T)$ is the superconducting energy gap at a temperature T , and k_B is Boltzmann's constant. For temperatures much lower than the critical temperature ($T \rightarrow 0$), this expression simplifies to

$$I_c R_N = \frac{\pi \Delta}{2e}. \quad (2.2)$$

The Josephson inductance L_J , which characterizes the nonlinear current-phase relationship of the junction, is given by

$$L_J = \frac{\Phi_0}{2\pi I_c} \quad (2.3)$$

where $\Phi_0 = h/(2e)$ is the magnetic flux quantum. Substituting Eq. (2.2) into this relation yields

$$L_J = \frac{\hbar R_N}{\pi \Delta} \quad (2.4)$$

This result shows that, in the ideal case, the Josephson inductance scales linearly with the normal-state resistance. In practice, deviations from the ideal Ambegaokar–Baratoff relation are frequently observed. To account for these effects, an empirical correction factor α is often introduced:

$$L_J \approx \alpha \frac{\hbar R_N}{\pi \Delta} \quad (2.5)$$

where α typically ranges between 1.2 and 1.4 for Al/AlO_x/Al Josephson junctions, providing better agreement between the theoretical prediction and experimental measurements.

The Josephson energy E_J , which represents the energy scale associated with tunneling Cooper pairs, is related to the critical current by

$$E_J = \frac{\Phi_0 I_c}{2\pi} = \frac{\pi \hbar \Delta}{4e^2 R_N} \quad (2.6)$$

This expression shows that E_J scales inversely with the normal-state resistance. When expressed in frequency units by dividing by Planck's constant, the relation becomes

$$\frac{E_J}{h} = \frac{\Delta}{8e^2 R_N} \quad (2.7)$$

and may likewise be corrected empirically as

$$\frac{E_J}{h} \approx \frac{1}{\alpha} \frac{\Delta}{8e^2 \alpha R_N} \quad (2.8)$$

Equations (2.5) and (2.8) provide a practical way to estimate the Josephson inductance and Josephson energy from simple room-temperature resistance measurements. By evaluating R_N for multiple junction sizes, we can predict how L_J and E_J scale with junction width, providing valuable design feedback prior to low-temperature characterization. In the following analysis, we use a superconducting gap of $\Delta_{\text{Al}} = 208 \mu\text{eV}$ for aluminum-based junctions [23].

Figure 2.16 shows the extracted parameters for a series of Manhattan junctions and SQUIDS with varying widths. Panel (a) shows the measured normal-state resistance R_N , which decreases with increasing junction width as expected from the larger effective area. Using the relations derived above, panels (b) and (c) show the corresponding Josephson inductance L_J and Josephson energy E_J . L_J and R_N have a linear dependence, while E_J follows the inverse trend. The data is fitted with a smooth polynomial to illustrate these scaling relations and to facilitate parameter estimation for device designs.

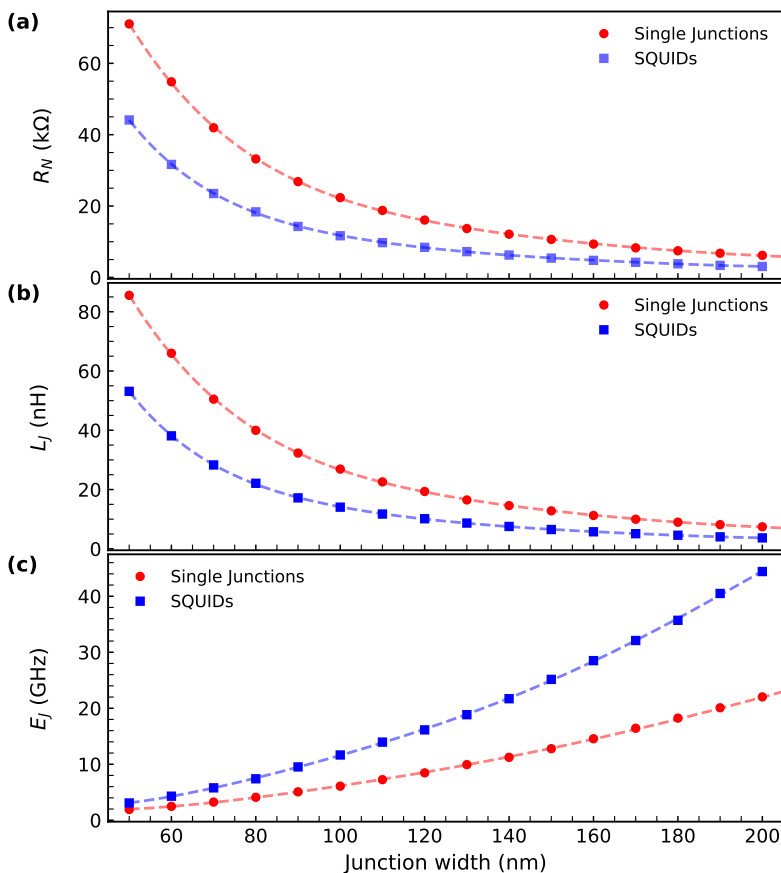


Figure 2.16: Measured Josephson junctions parameters as a function of junction width. (a) Room temperature measurements of the normal-state resistance R_N for junctions fabricated with a range of widths. **(b)** Josephson inductance L_J and **(c)** Josephson energy E_J calculated from R_N based on the Ambegaokar–Baratoff relation with an empirical correction factor $\alpha = 1.3$ to account for non-idealities.

2.2.3. DEVICE PACKAGING

After the chip fabrication is completed, the chips are integrated into a printed circuit board (PCB) that provides a thermal connection to the rest of the cryogenic setup, and also the electrical interface for the microwave signals. The PCB has a recessed pocket in which the chip is placed. This pocket ensures that the top of the chip lies on the same height as the PCB surface, which helps to keep the wire bond length as short as possible. The chip is glued to the PCB using GE varnish (GE 7031). A small amount of varnish is applied to each corner of the pocket before the chip is placed inside, where it is left to air dry for at least 30 minutes at room temperature. This fixes the chip in place, which is important for wire bonding, and it also acts as a thermal connection between the chip to the PCB. An overview of the PCB packaging and wire-bonding scheme is shown in Fig. 2.17.

Wirebonding is performed using a West-Bond 4KE semi-automatic wire bonder. Aluminium bonding wire is used, as it becomes superconducting at cryogenic temperatures and therefore minimizes adding additional resistive losses in the signal path. Typically, two or three parallel wirebonds are used to connect the on-chip launch pads to the corresponding signal traces on the PCB. In addition, on chip wirebonds are used to suppress parasitic slotline modes, which can lead to microwave crosstalk and reduced coherence. Finally, many grounding wirebonds are added along the entire perimeter of the chip. This creates the ground connection between the chip and the PCB. Furthermore, it also helps with the thermalization of the chip during the cooldown. We also transitioned to a new PCB design that incorporates Rosenberger spring-loaded coax connectors. This approach replaced the previous method used in our group that relied on soldered SMP connectors for every new board. The spring-loaded coax connector simplifies the assembly process. Furthermore, this design change also allowed us to reduce the size of the PCB and the rest of our sample packaging. Fig. 2.17 e shows the size comparison between the previous PCB on the left ($\varnothing$ 48 mm) and the new design on the right ($\varnothing$ 30 mm).

Figure 2.18 shows how the PCB is integrated into a copper sample box for cryogenic measurements. The PCB is positioned inside the sample box using three alignment holes that mate with corresponding alignment pins in the copper package. This defines a well-controlled orientation and ensures a reproducible position of the PCB to the sample box and the coaxial connectors.

The sample box has a two-part copper lid. The first lid section is placed directly on top of the PCB and is screwed into the bottom part of the box, forming a light-tight enclosure around the device. This also defines the positions of the RF connector feedthroughs, which are directly positioned above the corresponding connector traces on the PCB. The spring-loaded RF connectors are then inserted through the corresponding holes in this first lid. A second copper lid is mounted on top of the first lid to mechanically compress an external spring that surrounds each connector. This external spring pushes the connector body downward onto the PCB, ensuring a reliable electrical contact between the outer conductor of the connector and the PCB ground pads. In addition, each connector contains an internal spring-loaded center pin, which independently ensures that there is good electrical contact between the center conductor and the corresponding signal pad on the PCB.

After assembly, the sample box is mounted onto a cold finger, onto which up to two sample boxes can be installed simultaneously, as shown in Fig. 2.18 b. The cold finger provides a mechanical and thermal connection between the sample box and the mixing chamber of the dilution refrigerator. When required, a coil magnet can be mounted directly underneath the sample box to enable the application of a local magnetic field at the device. To suppress external magnetic flux noise, the cold finger and sample boxes are enclosed by an inner aluminum shield and an outer Cryoperm magnetic shield. Small openings in the shielding lids allow the RF coaxial cables to exit the enclosure. The complete assembly is then mounted to the underside of the mixing chamber plate.

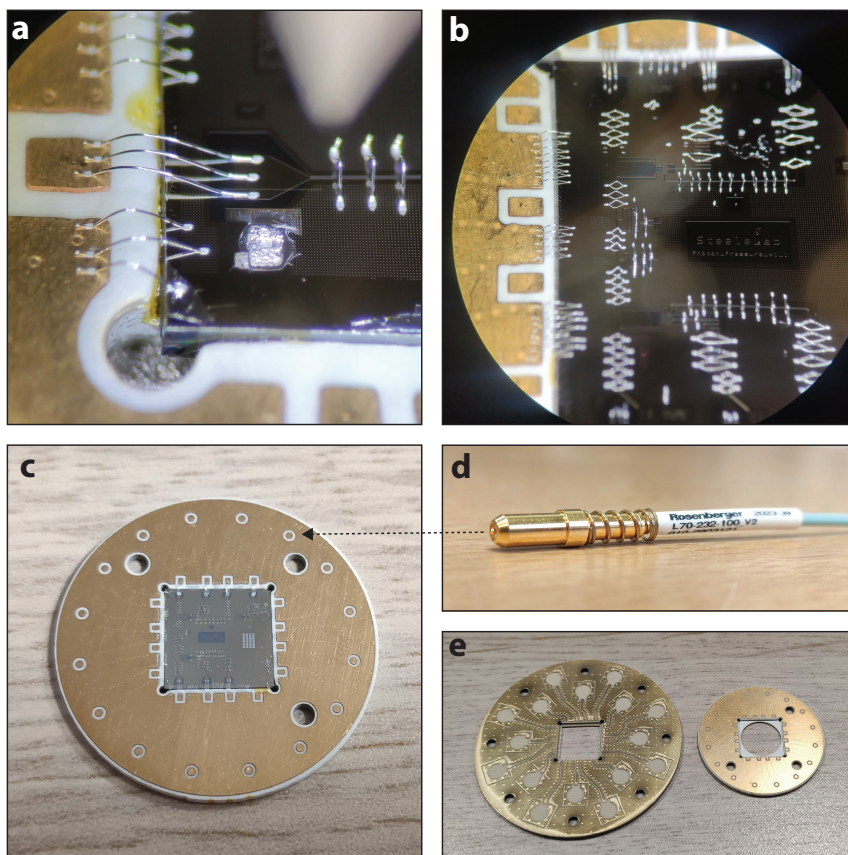


Figure 2.17: Wirebonds and PCB packaging. (a) Chip mounted inside the recessed pocket of the PCB. Excess GE varnish used to glue the chip at the four corners can be seen at the edge in yellow. Wirebonds connect the on-chip-launchpads to the corresponding signal traces on the PCB. (b) Overview of the same device showing the on-chip and grounding wirebonds. (c) Overview of the Entire PCB with the chip mounted and wirebonded. (d) Photograph of the spring-loaded coax connector. When placed in the sample box, the large spring presses the connector against the PCB, while the center conductor pin is also spring-loaded internally. (e) Size comparison between the previous PCB design (left, $\varnothing$ 48 mm) and the new PCB design (right, $\varnothing$ 30 mm).

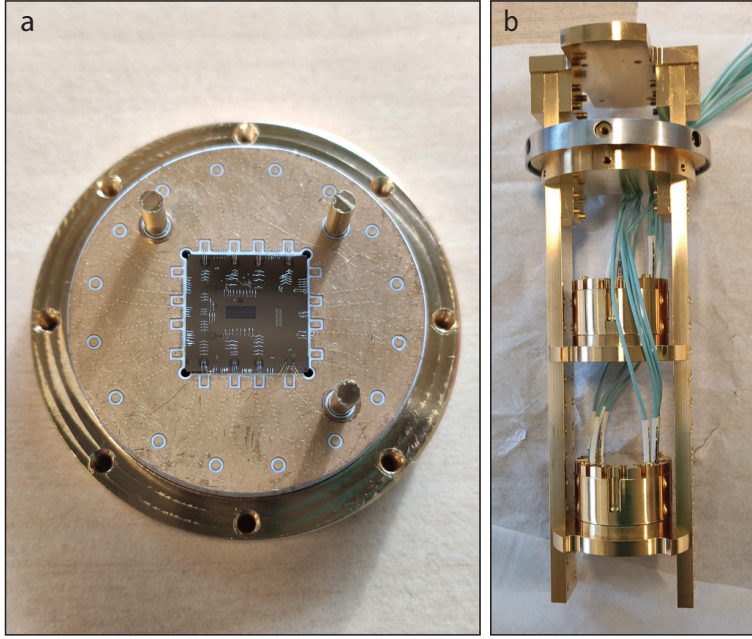


Figure 2.18: Sample mounting and packaging. (a) PCB placed on the bottom part of the copper sample box. Three alignment pins are used to accurately position the PCB within the package and ensure a reproducible chip orientation relative to the package. (b) Two assembled sample boxes mounted onto the cold finger, which is later screwed into the bottom plate of the dilution refrigerator mixing chamber. Prior to mounting, the samples are enclosed in an aluminium shield and an outer Cryoperm magnetic shield to reduce magnetic flux noise.

2.3. CRYOGENIC MEASUREMENT SETUP

To observe quantum mechanical behavior in superconducting circuits, the devices must be operated at sufficiently low temperatures to suppress thermal excitations. This requires that the thermal energy $k_B T$ is much smaller than the energy spacing of the system. For a circuit with transition frequency ω , this condition can be expressed as

$$k_B T \ll \hbar \omega. \quad (2.9)$$

Superconducting microwave resonators and qubits typically operate at frequencies in the range of a few gigahertz. For a circuit at 5 GHz, this corresponds to an equivalent temperature of approximately 240 mK. To strongly suppress thermal excitations, the operating temperature must therefore be well below this value. Such temperatures cannot be reached using conventional cryogenic techniques such as cooling in a liquid helium bath. Immersion in liquid ^4He cools a sample to 4.2 K, and pumping on the helium bath can further reduce the temperature to approximately 1.2 K. These temperatures remain far above the relevant energy scales of superconducting quantum circuits.

Temperatures on the order of 10–25 mK are obtained using a $^3\text{He}/^4\text{He}$ dilution refrigerator [24]. In these systems, cooling relies on the process by which ^3He atoms mix with ^4He at very low temperatures. Below a critical temperature, a mixture of the two helium isotopes separates into two distinct phases. A concentrated phase that is rich in ^3He , and a dilute phase in which a small fraction of ^3He is dissolved in ^4He . These two phases coexist at the mixing chamber.

In a cryogen-free dilution refrigerator, the initial cooling to cryogenic temperatures is provided by a pulse-tube cryocooler, which is a closed system that cools the system to approximately 4 K without the use of a liquid helium bath. Further cooling towards the millikelvin regime is achieved through the successive cooling stages of the dilution refrigerator, with the lowest temperatures reached at the mixing chamber. At this stage, when ^3He atoms move from the concentrated phase into the dilute phase, energy is required for the atoms to dissolve into the dilute mixture. This energy is taken from the surrounding environment, which extracts heat from the mixing chamber. As a result, the mixing chamber plate and the experiments thermally anchored to it are cooled.

Continuous cooling is maintained by circulating ^3He through the system. Pumping on the dilute phase at the still removes ^3He atoms, which drives a steady flow of ^3He across the phase boundary at the mixing chamber. While the available cooling power at the mixing chamber is limited, typically on the order of tens of microwatts, it is sufficient to maintain stable millikelvin temperatures when heat loads are properly managed in the experiments [25].

The cryogenic measurements presented in this thesis were performed in a cryogen-free dilution refrigerator (BlueFors LD250) operating at millikelvin temperatures. A photograph of the system is shown in Fig. 2.19. The samples are mounted at the bottom of the mixing chamber plate. Microwave signals are routed between room temperature and the mixing chamber plate using coaxial cables chosen for their low thermal conductivity and low signal loss. Stainless-steel coaxial cables are used from room temperature down to the 4 K stage. Below 4 K, superconducting NbTi coaxial cables are used. To suppress thermal noise and to ensure proper thermalization of the microwave lines, cryogenic attenuators are distributed along the input lines at multiple temperature stages. On the output lines, cryogenic microwave components such as isolators and circulators are used to prevent noise generated at the higher-temperature stages of the measurement chain from propagating back toward the device. High-electron-mobility transistor (HEMT) amplifiers mounted at the 4 K stage typically provide the first stage of amplification for the weak microwave signals before further amplification at room temperature.

The precise wiring configuration, including the number and placement of attenuators, filters, and amplifiers, depends on the specific experiment and device under investigation. Detailed wiring diagrams and component values are therefore provided in the Supplementary Information accompanying each research chapter.

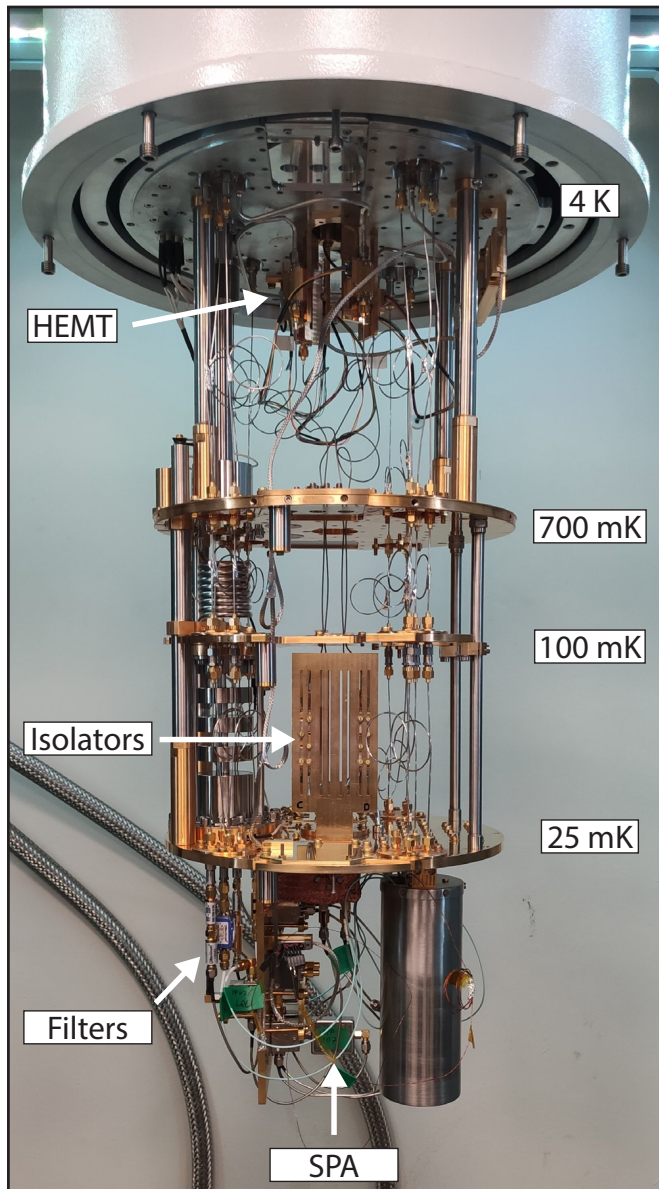


Figure 2.19: BlueFors LD250 dilution refrigerator. Photograph of the inside of a BlueFors LD250 dilution refrigerator, showing the cold finger (enclosed in magnetic shielding) mounted beneath the mixing-chamber plate. The labeled temperature stages (4 K, 700 mK, 100 mK, and 25 mK) and the locations of selected components of the microwave measurement chain are indicated, including the 4 K HEMT amplifier, isolators, filters, and a SNAIL parametric amplifier (SPA).

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3

STRONG INTRINSIC LONGITUDINAL COUPLING IN CIRCUIT QUANTUM ELECTRODYNAMICS

Radiation-pressure interactions between harmonic oscillators have enabled exquisite measurement precision and control, made possible by using strong sideband drives, enhancing the coupling rate while also linearizing the interaction. In this Letter, we demonstrate a strong intrinsic longitudinal coupling, a circuit quantum electrodynamics analog of the radiation-pressure interaction, between a transmon qubit and a linear microwave resonator. A red-detuned sideband drive results in an on-demand Jaynes-Cummings interaction with a high on-off ratio. We measure a longitudinal coupling rate an order of magnitude larger than all decay rates, placing the device in the strong coupling regime. The intrinsic longitudinal interaction demonstrated here will enable the development of high-connectivity quantum information processing hardware and the exploration of the gravitational decoherence of quantum objects.

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3.1. INTRODUCTION

Radiation-pressure interactions arise when the frequency of a system is modulated by the coordinate of a second system. Interactions of this type between harmonic oscillators have provided the foundation for a wide range of groundbreaking experiments in the last several decades. A canonical example is that of cavity optomechanics [1] but includes related fields such as cavity magnomechanics [2–5], photon-pressure coupling [6–10], and the center of mass motion of trapped ions [11, 12]. Such interactions have enabled scientific achievements such as ground-state cooling of mechanical objects [13–15], the generation of deterministic mechanical entanglement [16], microwave-to-optical frequency conversion [17, 18], and quantum non-demolition measurements [19, 20]. These demonstrations have relied on enhancing the typically low coupling rate via strong sideband drive tones, linearizing the intrinsically non-linear interaction.

In the context of superconducting qubits, the interaction analogous to radiation-pressure known as "longitudinal coupling" has recently been proposed as a new paradigm for quantum information processing [21, 22] with exciting potential applications for quantum processors [23, 24], and has been widely demonstrated in ion-trap hardware [11, 12, 25]. For example, longitudinal coupling could allow for all-to-all qubit connectivity and pure quantum non-demolition qubit readout [26–28]. The application of coherent drives has enabled an effective longitudinal coupling [29], which has been used to demonstrate efficient qubit readout [30, 31]. However, despite this interest, limited experimental progress on pure longitudinal coupling in superconducting circuits has been made, with most experimental realizations suffering from dominant parasitic transverse coupling spoiling the pure longitudinal nature [32].

In this Letter, we demonstrate an entirely hardware-realized *pure longitudinal* interaction between a superconducting qubit and a linear microwave resonator circuit. The device implements longitudinal coupling with negligible transverse coupling encoded directly in the circuit design. The interaction between the qubit and the resonator is engineered by threading the magnetic flux generated by the linear resonator through a superconducting quantum interference device (SQUID) [6, 7, 22], forming the inductance of the qubit. Using a sideband drive, we induce an effective Jaynes-Cummings interaction in the strong coupling limit that can be dynamically controlled with a high on-off ratio. Furthermore, the ac-Stark shift of the qubit provides an absolute photon number calibration [33], allowing for the extraction of the longitudinal coupling rate. We find a longitudinal coupling rate, $g_0 = 2\pi \times 11.9$ MHz, 2 orders of magnitude larger than the largest dissipation rate of the system, placing our system deep in the single-photon strong coupling regime [1].

The single-photon strong coupling regime has long been sought as it fully utilizes the inherent nonlinearity of the radiation-pressure interaction [34, 35]. Since linear interactions can only perform bilinear operations on Gaussian states, single-photon strong coupling would allow the dissipative engineering of states with Wigner negativity, such as described in the protocol for the generation of Schrödinger cat states [36]. Moreover, the longitudinal interaction allows a large number of photons to be driven into the linear resonator, unlike in conventional transverse circuit quantum electrodynamics (cQED) [37]. Therefore, combining longitudinal cQED with conventional optomechanics integrated within the linear resonator will enable the generation of mechanical Schrödinger cat states, via a double state swap procedure, to explore gravity's effect on quantum mechanics [38].

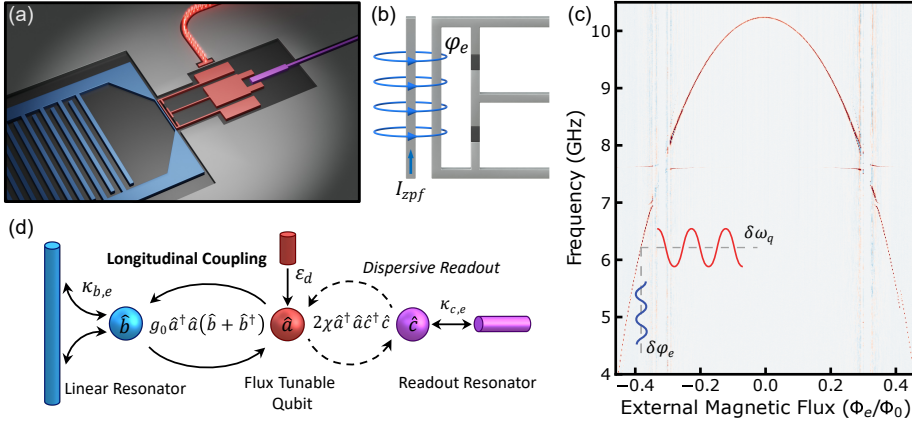


Figure 3.1: A superconducting quantum circuit with pure longitudinal coupling. (a) A rendering of the superconducting device. (b) Schematic depiction of the longitudinal coupling mechanism. Zero-point current in the linear resonator generates a magnetic flux that modulates the qubit frequency. (c) Qubit two-tone spectroscopy as a function of the external magnetic flux. The qubit has a sweet spot frequency of 10.2 GHz. The magnetic flux generated by currents in the linear resonator modulates the qubit frequency, resulting in longitudinal coupling. Data near the readout resonator frequency of 7.6 GHz is excluded due to the qubit resonator hybridization. (d) Schematic diagram of the device. The qubit $\hat{a}$ is coupled via a longitudinal interaction to the linear resonator $\hat{b}$ and coupled capacitively to a CPW resonator $\hat{c}$.

3.2. DEVICE DESIGN

The device consists of a modified transmon qubit coupled to two superconducting resonators, all with frequencies in the gigahertz range. The device's full circuit schematic can be seen in Fig. 3.1. All circuits are fabricated from niobium titanium nitride [39]; the fabrication details are outlined in supplementary information 3.7. A $\lambda/2$ coplanar waveguide resonator (CPW) labeled $\hat{c}$ is capacitively coupled to the qubit and acts as a standard readout resonator for qubit measurements and state preparation [37, 40]. It has a frequency $\omega_c = 2\pi \times 7.612$ GHz, a qubit-photon coupling rate $g_{ac} = 2\pi \times 118$ MHz, an internal and external coupling rate of $\kappa_{\text{int},c} = 2\pi \times 72.8$ kHz and $\kappa_{\text{ext},c} = 2\pi \times 673.8$ kHz.

The linear resonator labeled $\hat{b}$ consists of a large interdigitated capacitor shunted with a thin inductor wire. The linear resonator has a resonance frequency $\omega_b = 2\pi \times 4.347$ GHz, an internal and external coupling rate of $\kappa_{\text{int},b} = 2\pi \times 28.0$ kHz and $\kappa_{\text{ext},b} = 2\pi \times 88.6$ kHz and a residual parasitic dipolar coupling to the qubit $g_{ab} = 2\pi \times 1.4$ MHz. The parasitic capacitive coupling was minimized by careful qubit design and results in a critical photon number at the qubit operational point of $n_{\text{crit}} \equiv \Delta^2/4g_{ab}^2 \approx 350\,000$, where Δ is the qubit-linear resonator detuning [37]. The reduced hybridization between the qubit and linear resonator is an essential feature of the qubit design; the low coupling rate g_{ab} allows relatively large coherent drives to be applied to the linear resonator. We observe a bare-to-dressed transition of the linear resonator with a steady state drive $|\langle \hat{b} \rangle|^2 \approx 20\,000$, which is approximately an order of magnitude smaller than the critical photon number. However, we expect the number of

drive photons to increase significantly even with moderate reductions of the parasitic dipolar coupling [41]. These large coherent drives will be necessary if the linear resonator is replaced with a cavity optomechanical device since large coherent tones are required for mechanical state manipulation [14].

Finally, the qubit labeled $\hat{a}$ is a modified pocket-style transmon qubit with Manhattan-style aluminum aluminum-oxide Josephson junctions [42]. As seen in Fig. 3.1, one of the capacitor pads of the transmon is split and wrapped around the other, and the SQUID loop is extended away from the capacitor pads. The SQUID loop is placed $\sim 1 \mu\text{m}$ from the thin inductor wire of the linear circuit. This design was implemented to reduce the electric dipole moment of the qubit, minimizing the parasitic dipole coupling to the linear resonator, as discussed previously. We believe the parasitic coupling likely results from junction asymmetry resulting in finite current through the linear inductor of the SQUID and charging of the surrounding ground plane.

A coil mounted below the sample provides an external magnetic field that can tune the frequency of the qubit; see Fig. 3.1(c). During this work, the qubit was operated at a frequency $\omega_q(\Phi_e) \sim 2\pi \times 6.10 \text{ GHz}$, where Φ_e is the static external magnetic flux thread through the SQUID. The qubit can be driven via a direct XY drive line, which was designed to be weakly coupled such that the Purcell limited linewidth κ_{pur} is much less than the qubit linewidth γ_q . Finally, we measured a T_2^* limited linewidth $\gamma_q = 2\pi \times 677 \text{ kHz}$; qubit characterization is described within the supplementary information. The chip is mounted on the baseplate of a commercial dilution refrigerator operating with a base temperature of $T \sim 15 \text{ mK}$. For details about the measurement and setup, see supplementary information 3.7.

3.3. OPERATIONAL PRINCIPLE

We model the qubit as a Kerr oscillator with a large anharmonicity [37], the bare Hamiltonian of the system can be written in the form,

$$\hat{\mathcal{H}}_0/\hbar = \omega_q(\Phi_e)\hat{a}^\dagger\hat{a} + \frac{\alpha}{2}\hat{a}^\dagger\hat{a}^\dagger\hat{a}\hat{a} + \omega_b\hat{b}^\dagger\hat{b}. \quad (3.1)$$

Here, $\alpha = -388 \text{ MHz}$ is the qubit anharmonicity at the operation point, $\hat{a}^{(\dagger)}$ and $\hat{b}^{(\dagger)}$ are the qubit and linear resonator annihilation (creation) operators, respectively.

Current flowing through the inductor wire of the linear resonator generates an oscillating magnetic flux through the SQUID loop, modulating the qubit's frequency and inducing the coupling between the qubit and linear resonator, see Figs. 3.1(b), (c). It can be shown, after some algebraic manipulation and keeping only first-order terms (see the supplementary information) that the interaction can be written in the simple form [43]

$$\hat{\mathcal{H}}_{\text{int}}/\hbar = g_0\hat{a}^\dagger\hat{a}(\hat{b} + \hat{b}^\dagger). \quad (3.2)$$

Here, g_0 is the longitudinal coupling rate and is given by

$$g_0 = \frac{\partial\omega_q}{\partial\Phi}\Phi_{\text{zpf}}, \quad (3.3)$$

where $\partial\omega_q/\partial\Phi$ is the qubit's flux sensitivity at the operational point, and Φ_{zpf} is the magnetic flux thread through the SQUID loop generated by the zero-point current fluctuations of the

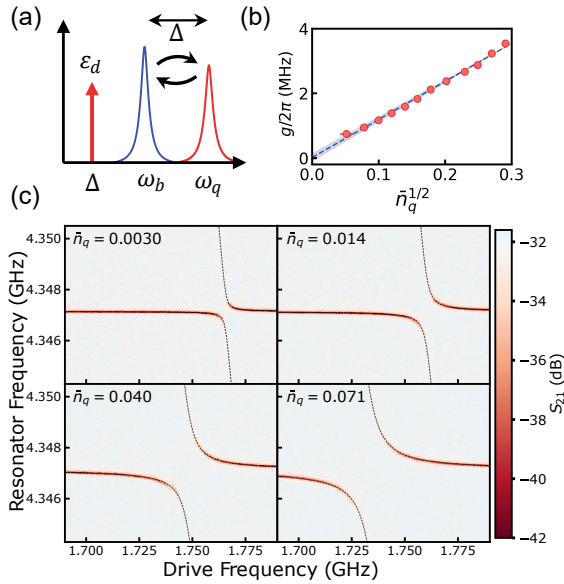


Figure 3.2: On-demand strong coupling by driving the longitudinal interaction. (a) Schematic diagram of the relative frequencies of the qubit (red), the linear resonator (blue), and the sideband drive (red). (b) Cavity-enhanced coupling rates extracted from the normal mode spectra at different sideband drive powers. The coupling rates are plotted as a function of the square root sideband photon number. The red points are data, and the blue dashed line is a linear fit to the data. The blue shaded area accounts for uncertainty in the bare qubit frequency. The photon number was calibrated using the ac-Stark shift. (c) The measured avoided level crossings of the linear resonator as a function of the sideband drive frequency. Starting from the top left, the drive powers considered here were $P_{sb} = 3.3, 8.3, 12.3, 14.4$ dBm, set at the source. The calibrated average number of sideband photons is printed in the top left of each panel. The dashed line is a fit; see the supplementary information for details.

linear resonator. The zero-point flux can be estimated from the mutual inductance and the zero-point fluctuations in the current given by $I_{zpf} = \sqrt{\hbar\omega_b/2L_b} \approx 36.8$ nA; with $\Phi_{zpf} = MI_{zpf}$. The mutual inductance can be estimated via the Biot-Savart law and gives $M \approx 25.9$ pH; therefore, we estimate the zero-point magnetic flux to have a value $\Phi_{zpf} \approx 461 \mu\Phi_0$. We have normalized the magnetic flux in units of $\Phi_0 = h/(2e)$. At the qubit operational point, the flux sensitivity $\partial\omega_q/\partial\Phi = 2\pi \times 26.0$ GHz/ Φ_0 , thus we can estimate the longitudinal coupling rate to have a value $g_0 \approx 2\pi \times 12.0$ MHz.

3.4. ON-DEMAND JAYNES-CUMMINGS

The on-demand drive-induced Jaynes-Cummings interaction was implemented using a sideband drive tone tuned to the difference frequency between the two coupled oscillators; see Fig. 3.2(a). In cavity optomechanics, the sideband drive is typically nearly resonant with the optical cavity, as mechanical frequencies are 1-10 MHz for most electromechanical systems [13, 15]. However, this device operates in a vastly different parameter regime. The sideband

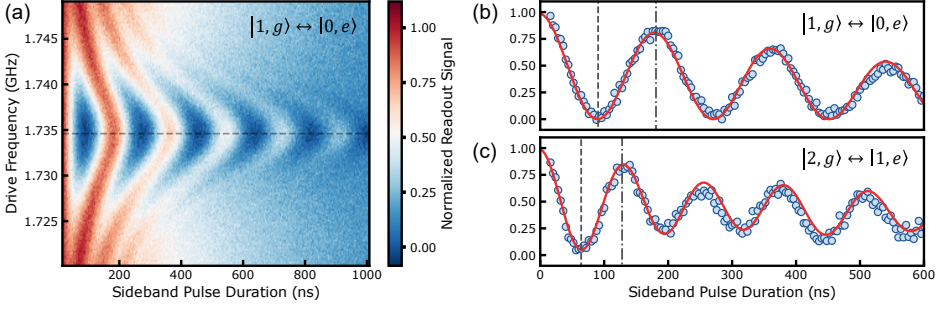


Figure 3.3: Observation of coherent swaps of photons up the on-demand Jaynes-Cummings ladder (a) Rabi chevrons with a single excitation in the system. The normalized output signal is proportional to the excited state probability of the qubit $P_{|e\rangle}$. (b) Line cut of resonant Rabi oscillations for the $|1, g\rangle \leftrightarrow |0, e\rangle$ transition. A Rabi frequency $g = 2\pi \times 2.76$ MHz was used in the numerical simulation. The dashed lines indicate the longitudinal π -pulse duration $t_\pi = 90.5$ ns. The dotted dashed lines indicate the Rabi oscillation period, and the solid red line is a numerical simulation of the Rabi oscillations. (c) Line cut of resonant Rabi oscillations for the $|2, g\rangle \leftrightarrow |1, e\rangle$ transition. An increased Rabi frequency $g_2 \approx \sqrt{2}g$ is observed. The dashed lines indicate the predicted second manifold longitudinal π -pulse duration $t_{\pi,2} = 90.5/\sqrt{2}$ ns. The dotted dashed lines indicate the Rabi oscillation period, and the solid red line is a numerical simulation of the Rabi oscillations.

drive frequency is below both the qubit and linear resonator frequencies. We should emphasize that we are driving single-photon sideband transitions and *not* the two-photon sideband transitions used in conventional cQED [44].

We can linearize the Hamiltonian about the sideband drive tone with frequency $\omega_d \approx \omega_q - \omega_b$ tuned near the red sideband applied via the XY drive line to the qubit. In a displaced frame corotating with the drive, we can linearize by defining $\hat{a} = \langle \hat{a} \rangle + \delta \hat{a}$; doing so, we can define the enhanced coupling rate as $g = g_0 |\langle \hat{a} \rangle|$, and $\bar{n}_q = |\langle \hat{a} \rangle|^2$ is the average sideband steady-state population of the qubit. The interaction Hamiltonian given in Eq. 3.2 can be written, taking the two-level approximation of the transmon and applying the rotating wave approximation, as

$$\hat{\mathcal{H}}_{\text{int}}/\hbar = g(\hat{\sigma}_- \hat{b}^\dagger + \hat{\sigma}_+ \hat{b}), \quad (3.4)$$

where $\hat{\sigma}_\pm$ are the qubit raising and lowering operators. The linearization process results in an on-demand drive-induced Jaynes-Cummings interaction created by the sideband drive; see the supplementary information for a full derivation [43].

The interaction is experimentally investigated by sweeping a sideband drive near the red sideband and probing the linear resonator with less than a single photon on average to remain in the single-excitation manifold of the Jaynes-Cummings ladder [45]. We observe a series of avoided level crossings with a correspondingly increasing coupling rate as a function of increasing sideband power and a shift to lower frequencies, as shown in Fig. 3.2(c). We can extract the cavity-enhanced coupling rate g from the avoided level crossings. However, we require an absolute calibration of the photon number to determine the longitudinal coupling rate. Conventionally, this is achieved by varying the temperature of the dilution refrigerator and performing thermal calibration [13], or applying a known modulation and carefully mea-

suming all loss and gain within the system [46, 47].

Instead, we can use the ac-Stark shift of the qubit for direct photon number calibration. As we increase the sideband drive power, the qubit experiences an increasing ac-Stark shift. The magnitude of the Stark shift is directly related to the average number of sideband photons in the qubit. One must take care when interpreting the Stark shift, which is why we initially modeled the qubit as a Kerr oscillator. Since the sideband drive is far detuned, the counter-rotating terms, known as the Bloch-Siegert shift [33], and the qubit anharmonicity contribute significantly. Considering the steady state of the Kerr oscillator and the ac-Stark shift, the steady-state qubit population for each sideband power can be directly calculated.

The value of the longitudinal coupling rate can be determined by considering the cavity-enhanced coupling rate as a function of the square root of the number of steady-state sideband photons, as seen in Fig. 3.2(b). Importantly, the relationship is linear, suggesting our approximations are valid and hold even for steady-state populations below the single-photon level, and the data has a y intercept of zero, which is expected since the cavity-enhanced coupling should vanish for no sideband drive. The value of the longitudinal coupling rate can be determined from the slope, giving a value of $g_0 = 2\pi \times 11.9 \pm 0.3$ MHz, which is in excellent agreement with our theoretical prediction and well within the single-photon strong coupling regime.

3.5. RABI OSCILLATIONS

Finally, to demonstrate the coherent nature of the interaction, we measure vacuum Rabi oscillations between the qubit and the linear resonator. The qubit is first prepared into its excited state $|e\rangle$ using a 24 ns resonant π -pulse. A square wave sideband pulse is applied near the red sideband frequency for a variable time τ . Following the sideband pulse, the qubit state is measured using the capacitively coupled CPW resonator. This protocol is repeated for various sideband drive frequencies and pulse durations, resulting in the vacuum Rabi oscillations shown in Fig. 3.3. We observe a typical Rabi “chevron” pattern with coherent exchanges between the qubit and the linear resonator. A line cut is fit using a numerical simulation of the Rabi oscillations to confirm our linear model; see the supplementary information [43]. From the numerical simulation, we can extract the coupling rate $g/2\pi = 2.76 \pm 0.03$ MHz, [48] resulting in a sideband π -pulse duration of ~ 91 ns; see Fig. 3.3(b). This is in excellent agreement with the measured normal mode splitting at nominally the same sideband power, $g/2\pi = 2.81 \pm 0.08$ MHz. Finally, we observe transitions to higher excited manifolds of the Jaynes-Cummings ladder by preparing the linear oscillator in the Fock state $|1\rangle$ using a sideband π -pulse and measuring the Rabi oscillations. We observe a Rabi frequency with a value $g_2 \approx \sqrt{2}g$, as predicted by the Jaynes-Cummings model [48]; see Fig. 3.3(c).

3.6. OUTLOOK

This Letter presents a unique interface between a superconducting qubit and a linear microwave resonator. The interaction was confirmed by performing normal-mode spectroscopy of the hybrid modes induced via a red sideband drive and observing coherent Rabi oscillations. The value of the longitudinal coupling rate was determined by calibrating the sideband

drive using the ac-Stark shift. Our device is well within the single-photon strong coupling limit.

The innovative qubit design provides a new platform for testing theories demanding single-photon strong coupling, such as dissipative Schrödinger cat state generation within the microwave resonator [36]. Moreover, the large critical photon number, $n_{\text{crit}} \approx 350\,000$, and the ability to dynamically modulate the coupling should allow for future integration with optomechanics experiments. In such an experiment, the cat state generated within the linear resonator can be swapped into a mechanical oscillator using a strong optomechanical sideband pulse, generating a macroscopic superposition allowing for exotic tests of gravitational decoherence [38]. Such a swap to the mechanical resonator would *not* be possible using conventional Jaynes-Cummings coupling due to the large driving requirements on the linear microwave cavity.

Moreover, this platform may also be used for interesting on-chip quantum optics experiments; for example, at the top sweet spot to first order, the linear coupling should be suppressed, i.e., $g_0 \approx 0$; however, there will remain quadratic coupling of the form $\hat{\mathcal{H}}_q/\hbar = g_q \hat{a}^\dagger \hat{a} (\hat{b} + \hat{b}^\dagger)^2$. Quadratic coupling has been long sought and was an initial driving force for the membrane in the middle experiments [49]. By eliminating linear coupling, the quadratic coupling may allow quantum non-demolition measurement of the photon number within the linear resonator. This enables an on-chip cQED measurement of stochastic photon jumps as they decay from the linear resonator [50, 51].

Finally, the longitudinal coupling can be explored as a unique architecture for quantum information systems [21, 23, 24, 26]. Multiple qubits could be coupled to a single bus resonator, and by applying appropriate sideband pulses, sideband drive-mediated qubit-qubit gates may be possible. All qubits connected via the longitudinal interaction to a single bus resonator would allow a fully superconducting implementation of efficient quantum computation architectures currently studied using ion traps [52].

3.6.1. AUTHORS CONTRIBUTIONS:

C.A.P. performed theoretical modeling, performed data analysis, designed the qubit, performed daily supervision of measurements and coordination of the team, wrote the first draft of the manuscript with input from all authors, and incorporated feedback from authors into the final manuscript. R.C.D. designed and simulated the full device layout, developed the fabrication recipe and fabricated the device, performed spectroscopic measurements, contributed to data analysis, and contributed to figure design. S.D. contributed to the verification of theoretical models. S.D. and E.S. performed time-domain qubit measurements and analysis of those measurements. G.A.S. was responsible for conceiving the experiment, overall supervision of the project, and the acquisition of funding for the project. All authors contributed to the formulation of the manuscript storyline and the composition of the figures.

3.6.2. DATA AND MATERIALS AVAILABILITY:

The data that support the findings of this article are openly available [53].

3.7. SUPPLEMENTARY INFORMATION

3.7.1. MATERIALS AND METHODS

DEVICE FABRICATION

The device is fabricated on 10x10 mm 525 μm thick high-resistivity silicon chips with 200 nm of niobium titanium nitride (NbTiN). The NbTiN is deposited by the Dutch Institute for Space Research (SRON) following the procedure described in Ref. [39]. The bulk of the circuitry is fabricated by spinning a (positive) e-beam resist layer (AR-P 6200.18, 4000 rpm) and patterning the CPW, linear resonator, qubit islands and feedlines using electron-beam lithography. After development (Pentylacetate, O-xylene, IPA) and a post-exposure bake, the NbTiN is removed using an SF_6 reactive ion etching step, followed by an in-situ oxygen plasma descum. After stripping the resist, the Josephson junctions are fabricated by depositing an MMA/PMMA resist bi-layer and patterning the junctions using electron-beam lithography. The resist is developed using cold H_2O :IPA (1:3) and IPA. The chip is dipped in a buffered oxide etchant (BOE) for 1 min prior to loading into the Plassys deposition system. The junctions are evaporated using a double-angle deposition technique, forming a pair of Manhattan-style Josephson junctions. After liftoff in NMP and wire bonding, the chips are loaded into the dilution refrigerator.

MEASUREMENT SETUP

All measurements were performed on the base plate of a dilution refrigerator operating at a base temperature $T \sim 15$ mK. A complete schematic of the wiring within the fridge is shown in Fig. 3.4. The $\lambda/2$ CPW readout resonator is connected to its own feedline via a cryogenic circulator and measured in reflection. The linear resonator is coupled in a notch-type geometry and measured in transmission using a different output line. Both output lines are passed through a pair of low-temperature isolators and via superconducting coaxial cables to a pair of HEMT amplifiers and additional amplification at room temperature. The input lines connecting to the $\lambda/2$ CPW resonator, the qubit drive line, and the input line of the linear resonator all have 59 dB of cryogenic attenuation, with additional room-temperature attenuation. The device is mounted in a light-tight copper box on a copper bracket within a superconducting aluminum and Mu-metal magnetic shield. An external superconducting coil magnet is mounted underneath the copper sample box and inside the magnetic shielding. Spectroscopic measurements were performed using a four-port Keysight PNA for both the qubit and linear resonator, and the time domain measurements were performed using a Quantum Machines OPX, Octave, and Rohde & Schwarz SGS100A vector microwave generator.

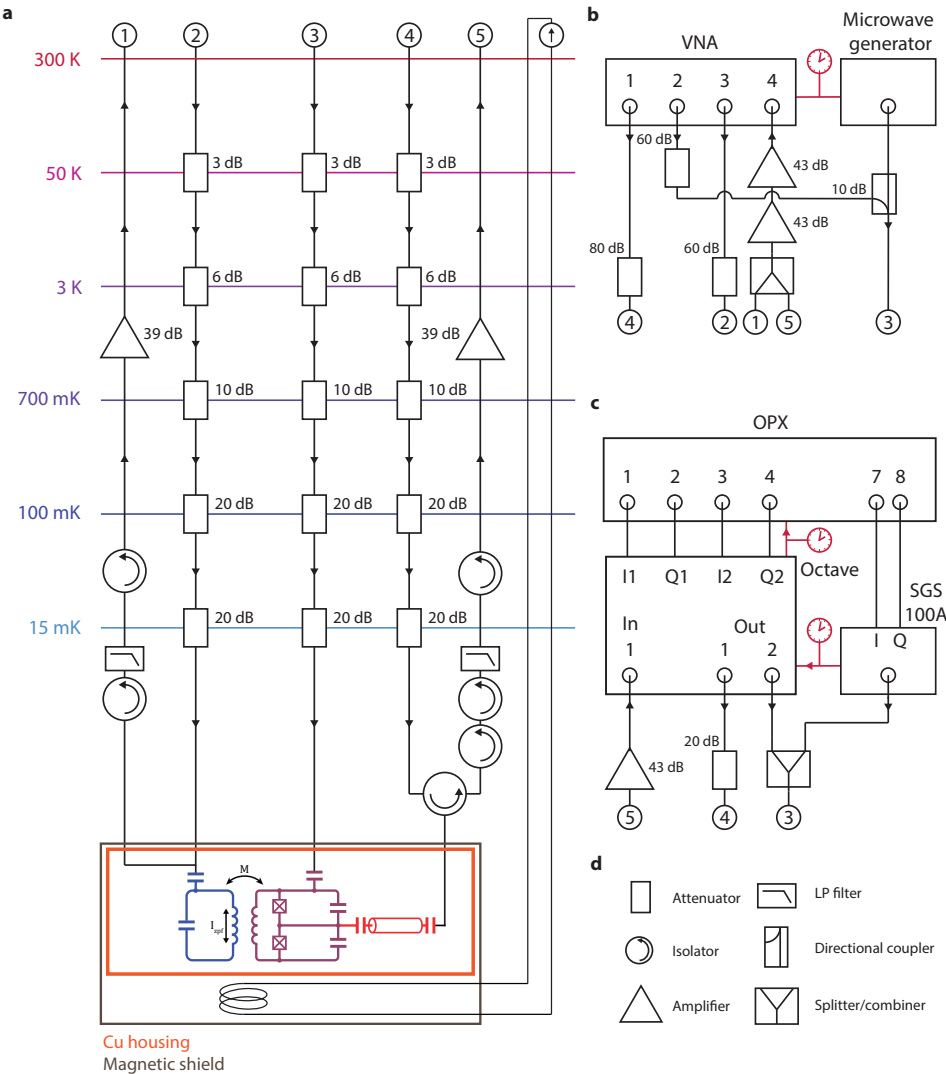


Figure 3.4: Schematic of the measurement setup. **a**, Wiring diagram inside the dilution refrigerator. **b**, Room temperature setup for VNA measurements. **c**, Setup for time domain measurements. **d**, Overview of components.

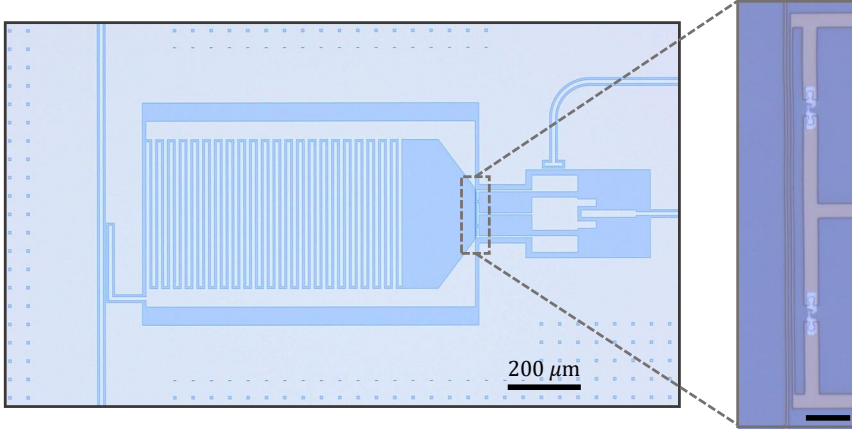


Figure 3.5: Superconducting circuit. a, Optical micrograph image of the qubit. The scale bar on the zoom is $15 \mu\text{m}$.

ZERO-POINT CURRENT FLUCTUATIONS

The linear RF circuit consists of a set of interdigitated capacitors (IDC) with a thin inductor wire; see Fig. 3.5. The capacitance of the IDC can be estimated analytically following Ref. [54],

$$C_{\text{IDC}} = (N - 3) \frac{C_1}{2} + 2 \frac{C_1 C_2}{C_1 + C_2} \quad (3.5)$$

where

$$C_i = 2\epsilon_0\epsilon_{\text{eff}}l \frac{K(k_i)}{K(k'_i)}, \quad i = 1, 2. \quad (3.6)$$

Here, $K(k_i)$ are the elliptic integrals of the first kind, l is the length of the IDC fingers, $\epsilon_{\text{eff}} = (\epsilon_r + 1)/2$ is the effective permittivity where the permittivity of silicon is $\epsilon_r = 11.8$, N is the total number of fingers and

$$k_1 = \sin\left(\frac{\pi}{2} \frac{a}{a+b}\right), \quad (3.7)$$

$$k_2 = 2 \frac{\sqrt{a(a+b)}}{2a+b}, \quad (3.8)$$

$$k'_i = \sqrt{1 - k_i^2}, \quad (3.9)$$

where, a is the finger width, and b is the gap between fingers. The fabricated linear resonator has $N = 44$, $a = 10\mu\text{m}$, $b = 6\mu\text{m}$, and $l = 400\mu\text{m}$. Therefore, Eq. 3.6 estimates the capacitance of the linear circuit as $C_{\text{IDC}} = 1.26 \text{ pF}$, which is in excellent agreement with the

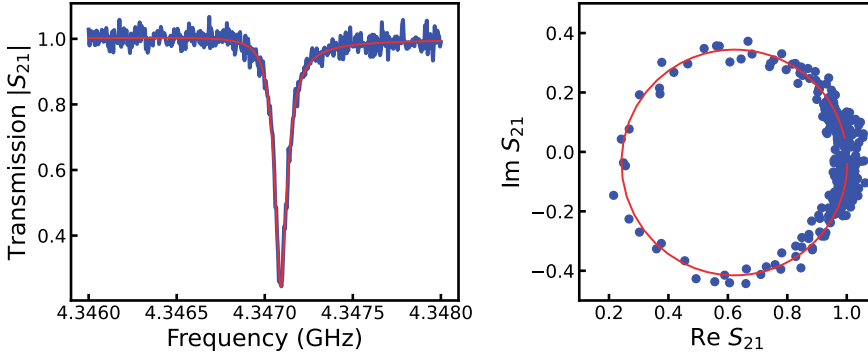


Figure 3.6: The linear circuit response. **a**, The magnitude of the transmission response $|S_{21}|$ of the linear resonator. **b**, The complex transmission data S_{21} near the resonance frequency of the linear resonator. In both plots, the circles are data; the line is a fit. The data plotted here has had its background subtracted.

Ansys capacitance simulation calculating a total capacitance of $C_b = 1.29$ pF, suggesting there is little stray capacitance from the remainder of the circuit. The measured resonance frequency of the circuit is $\omega_b = 2\pi \times 4.347$ GHz. Assuming the IDC capacitance is the only capacitance in the circuit, we can determine the circuit inductance by

$$\omega_b = \frac{1}{\sqrt{L_b C_b}}. \quad (3.10)$$

giving a total inductance of $L_b = 1.06$ nH. Therefore, the current zero-point fluctuations through the inductor of the linear circuit are given by

$$I_{\text{zpf}} = \sqrt{\frac{\hbar \omega_b}{2L_b}} \quad (3.11)$$

with the values calculated above, the zero-point current $I_{\text{zpf}} \approx 36.8$ nA.

3.7.2. DEVICE CHARACTERIZATION

LINEAR CIRCUIT PARAMETER EXTRACTION

The linear microwave resonator is coupled to a transmission line in the ‘notch’ style geometry. Thus, the data is fitted to the transmission response function

$$S_{21}^{\text{notch}}(\omega) = 1 - \frac{\kappa_{\text{ext}} e^{i\theta}}{(\kappa_{\text{int}} + \kappa_{\text{ext}}) + 2i(\omega - \omega_b)}. \quad (3.12)$$

Here, θ accounts for any impedance mismatch, ω_b is the resonance frequency, and κ_{int} and κ_{ext} are the internal decay rate and the decay rate into the transmission line, respectively. The fit of the measured transmission spectrum is shown in Fig. 3.6. We extract the parameters $\omega_b = 2\pi \times 4.347$ GHz, $\kappa_{\text{int}} = 2\pi \times 28.0$ kHz, $\kappa_{\text{ext}} = 2\pi \times 88.6$ kHz, and $\kappa = 2\pi \times 116.6$ kHz.

3.7.3. QUBIT CHARACTERIZATION

The initial characterization of the qubit was done using two-tone spectroscopy. A weak pump tone was applied near the qubit frequency while simultaneously monitoring the frequency of the readout CPW. This technique allows the determination of the qubit frequency and the anharmonicity of the qubit. By sweeping the external magnetic field, we are able to map out the qubit flux arc. We found a sweet spot frequency $\omega_q = 2\pi \times 10.2$ GHz and could tune the qubit frequency below 500 MHz. Next, with the qubit at its operational point, by increasing the qubit pump tone, we are able to measure the anharmonicity by driving the $|1\rangle \rightarrow |2\rangle$ transition. We find an anharmonicity $\alpha = -2\pi \times 388$ MHz.

Next, the qubit was characterized by performing time domain measurements of its decay characteristics. At the operational point, the qubit has a $T_1 = 664$ ns and a $T_2^* = 235$ ns. These values are not state-of-the-art; however, they are not a result of the geometry since standard Xmon-style qubits fabricated using the same process show similar values for the characteristic decay rates. We have identified our most likely limiting factor to be our SF_6 etch step damaging the silicon surface prior to our junction deposition. This is currently a work in progress, and we aim to improve our coherence time with future devices. Secondly, we expect flux noise to limit our T_2^* ; this can be improved in future experiments by reducing the SQUID area and optimizing the mutual inductance to maintain a large g_0 . The current SQUID loop has an area $\sim 360 \mu\text{m}^2$, likely resulting in large amounts of flux noise. The SQUID loop size can be reduced with device optimization, reducing flux noise.

3.7.4. QUBIT-RESONATOR COUPLING

PARASITIC COUPLING

As discussed in the main text, the qubit was designed to minimize dipole coupling between the qubit and the linear resonator. The linear interaction of the form $\mathcal{H}_{\text{JC}}/\hbar = g_{\text{ab}}(\hat{\sigma}_+ \hat{b} + \hat{\sigma}_- \hat{b}^\dagger)$ results in the hybridization of the qubit with the linear resonator. Even in the dispersive regime, where $\Delta \gg g_{\text{ab}}$, the weak hybridization results in a critical photon number, $n_{\text{crit}} = \Delta^2/(4g_{\text{ab}}^2)$, in which the dispersive approximation breaks down [40]. Here, $\hat{\sigma}_\pm$ are the qubit raising and lowering operators, $\hat{b}^{(\dagger)}$ are the linear circuit annihilation (creation) operators and $\Delta = |\omega_q - \omega_b|$ is the qubit detuning from the linear resonator.

Flux tuning the qubit to be approximately resonant with the linear resonator results in an avoided level crossing, as shown in Fig. 3.7. The dipole coupling is given by the frequency difference between the upper and lower normal modes when the qubit and linear circuit are resonant. We extract a dipole coupling rate of $g_{\text{ab}} \approx 2\pi \times 1.4$ MHz, several orders of magnitude smaller than current state-of-the-art circuit quantum electrodynamics devices [40]. The qubit operational point used in this work was detuned from the linear circuit by $\Delta = 1.75$ GHz; therefore, we estimate a critical photon number within the linear circuit of $n_{\text{crit}} \approx 350\,000$.

LONGITUDINAL COUPLING

Longitudinal coupling arises from a radiation-pressure-like interaction between the qubit and linear resonator. This coupling modulates the qubit frequency due to the mutual inductive coupling. We can consider the modulation of the qubit frequency by Taylor expanding about

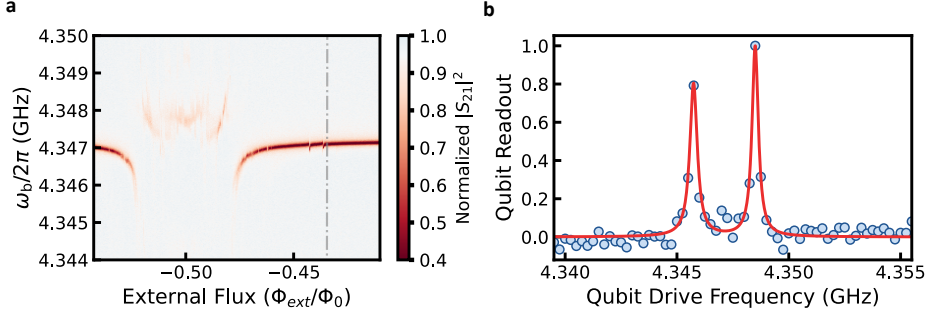


Figure 3.7: Avoided level crossing between the qubit and linear resonator **a**, Measured transmission spectra $|S_{21}|^2$ of the linear circuit as the qubit — shown by the grey dashed line — is swept through the linear resonator’s frequency. Additional avoided crossings are observed as the qubit tunes to its lower sweet spot; the origin of these additional splittings near a half flux quantum is currently being investigated. **b**, Two-tone spectroscopy of the qubit fully hybridized with the linear resonator. We can extract the linear coupling coefficient g_{ab} between the linear circuit and the qubit from the avoided crossing.

a set flux point, given by

$$\omega_q(\Phi) = \omega_{q,0} + \frac{\partial \omega_q}{\partial \Phi} \Phi_{\text{zpf}} (\hat{b} + \hat{b}^\dagger) + \frac{1}{2} \frac{\partial^2 \omega_q}{\partial \Phi^2} \Phi_{\text{zpf}}^2 (\hat{b} + \hat{b}^\dagger)^2 + \dots \quad (3.13)$$

Where $\hat{b}^{(\dagger)}$ is the creation (annihilation) operator for the linear circuit. Inserting Eq. 3.13 into the Hamiltonian $\hat{\mathcal{H}} = \hbar \omega_q \hat{a}^\dagger \hat{a}$, we arrive at

$$\hat{\mathcal{H}}/\hbar = \left(\omega_{q,0} + \frac{\partial \omega_q}{\partial \Phi} \Phi_{\text{zpf}} (\hat{b} + \hat{b}^\dagger) + \frac{1}{2} \frac{\partial^2 \omega_q}{\partial \Phi^2} \Phi_{\text{zpf}}^2 (\hat{b} + \hat{b}^\dagger)^2 + \dots \right) \hat{a}^\dagger \hat{a}. \quad (3.14)$$

We can rewrite this in the form

$$\hat{\mathcal{H}}/\hbar = \omega_{q,0} \hat{a}^\dagger \hat{a} + g_0 \hat{a}^\dagger \hat{a} (\hat{b} + \hat{b}^\dagger) + g_q \hat{a}^\dagger \hat{a} (\hat{b} + \hat{b}^\dagger)^2 + \dots \quad (3.15)$$

Where we can recognize the second term as the longitudinal interaction we are interested in and the third term as the quadratic longitudinal interaction.

COUPLING RATE ESTIMATION

The mutual inductive coupling between the inductor wire of the linear circuit and the SQUID loop of the qubit quantifies the coupling between the linear circuit and the qubit. As shown in Fig. 3.5, the SQUID loop of the qubit has a designed length $l = 120\mu\text{m}$ and width $w = 3\mu\text{m}$, giving a nominal SQUID area of $360\mu\text{m}^2$; however, the final dimensions were slightly larger due to over-etching. The $1.0\mu\text{m}$ wide inductor wire passes $\sim 1.0\mu\text{m}$ from the SQUID loop, with measured geometric distances $d_1 \approx 1.8\mu\text{m}$ and $d_2 \approx 5.3\mu\text{m}$ are the smallest and largest distance between the inductor and the SQUID loop, respectively; see Fig. 3.5. The mutual

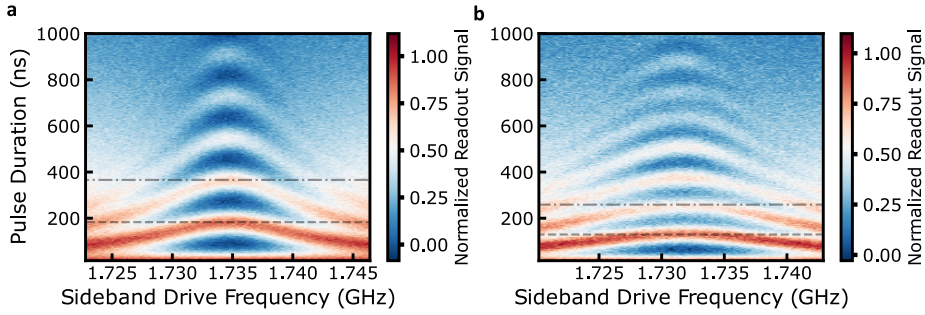


Figure 3.8: On-demand Jaynes-Cummings Ladder **a**, Rabi chevrons with a single excitation in the system. A Rabi frequency $g = 2\pi \times 2.73$ MHz was measured. **b**, Rabi chevrons with two total excitations in the system. The linear resonator is prepared into the Fock state $|1\rangle$ using a π -pulse. Following the state preparation, the Rabi chevron was measured. An increased Rabi frequency $g_2 \approx \sqrt{2}g$. The dashed lines indicate the Rabi oscillation period, and the dotted dashed lines indicate two Rabi oscillation periods.

inductance can be calculated by integrating the magnetic flux generated by the inductor wire through the SQUID loop and is given by

$$M = \frac{\mu_0 I}{2\pi} \ln \left(\frac{d_2}{d_1} \right). \quad (3.16)$$

Using this approximation, we find a mutual inductance $M \approx 25.9$ pH. Moreover, we can determine the zero-point fluctuations of the magnetic flux using the zero-point current, giving $\Phi_{zpf} = M I_{zpf} \approx 461 \mu\Phi_0$ where we have written the zero-point magnetic-flux fluctuations in units of flux quanta $\Phi_0 = h/(2e)$.

The longitudinal coupling rate is defined as

$$g_0 = \frac{\partial \omega_q}{\partial \Phi} \Phi_{zpf}. \quad (3.17)$$

The longitudinal coupling rate is the frequency shift of the qubit due to the zero-point fluctuations of the magnetic flux generated by the linear circuit. The frequency responsivity of the qubit can be determined by sweeping the external magnetic field at the operation frequency; we find $|\partial \omega_q / \partial \Phi| \approx 2\pi \times 26.0$ GHz/ Φ_0 . Therefore, at the operation point, we estimate a longitudinal coupling rate of $g_0 \approx 2\pi \times 12.0$ MHz, which is in excellent agreement with the experimentally measured value.

QUADRATIC COUPLING

If we consider the second-order Taylor expansion of the coupling, we can define the quadratic coupling between the qubit and linear resonator as

$$\hat{\mathcal{H}}_q / \hbar = g_q \hat{a}^\dagger \hat{a} (\hat{b} + \hat{b}^\dagger)^2, \quad (3.18)$$

Where the quadratic coupling rate is defined as,

$$g_q = \frac{1}{2} \frac{\partial^2 \omega_q}{\partial \Phi^2} \Phi_{zpf}^2. \quad (3.19)$$

Interestingly, the quadratic coupling is non-zero at the qubit sweet spot, where the linear coupling is precisely equal to zero. From the flux arc, at the sweet spot and the qubit operation point, we can estimate the quadratic coupling rate as $g_q \approx 2\pi \times 5.5$ kHz, and $g_q \approx 2\pi \times 14.0$ kHz, respectively. Keeping only the non-rotating terms, this quadratic coupling results in a cross-Kerr style coupling of the form,

$$\hat{\mathcal{H}}_q/\hbar \rightarrow \chi_q \hat{a}^\dagger \hat{a} \hat{b}^\dagger \hat{b}. \quad (3.20)$$

Where the value of the cross-Kerr coupling is $\chi_q = 2g_q \approx 28.0$ kHz.

It should be noted that the bare-to-dressed transition discussed in Section 3.7.4 is a result of the breakdown of a dispersive approximation. The non-perturbative cross-Kerr resulting from the quadratic interaction should not result in a critical photon number [27]. However, the non-perturbative cross-Kerr coupling may play a role in qubit dephasing.

ON-DEMAND JAYNES-CUMMINGS

To confirm that the interaction is indeed described by the Jaynes-Cummings Hamiltonian, we compared the single and two-excitation Rabi frequencies. The single excitation Rabi frequency was measured as described in the main text. The data was fit using a numerical simulation described below, where we excite the qubit into the state $|1\rangle$ using a resonant π -pulse. To determine the two-excitation Rabi frequency, we first initialize the linear resonator into the Fock state $|1\rangle$ by performing a sideband π pulse. This was followed by a measurement of the Rabi chevron, shown in Fig. 3.9(b). We observe a clear increase in the Rabi frequency as predicted by the Jaynes-Cummings model; the Rabi frequency increases by a factor of approximately square root two. We use our numerical simulation to determine the Rabi frequency given by $g = 2\pi \times 2.76$ MHz, which is in good agreement with a fit of the line-cut to a decaying sinusoidal function giving a value $g = 2\pi \times 2.73$ MHz. We attribute any small discrepancy to the fidelity of our initial preparation of the linear resonator Fock state. Due to losses within the system, on average, the state of the linear resonator is likely a mixed state of $|0\rangle$ and $|1\rangle$, reducing the ensemble-averaged Rabi frequency measured here. However, this measurement does confirm that the interaction produces a Jaynes-Cummings ladder. With device optimization, this will be explored in future work.

MASTER EQUATION SIMULATION

To model the dynamics of the drive-induced Jaynes-Cummings interaction, we simulate the system using the Lindblad Master equation using the Python package Qutip [55]. The dynamics are governed by the equation given by,

$$\frac{\partial \rho}{\partial t} = -\frac{i}{\hbar} [\hat{\mathcal{H}}(t), \rho] + \kappa_b \mathcal{L}[\hat{b}] \rho + \gamma_1 \mathcal{L}[\hat{\sigma}_-] \rho, \quad (3.21)$$

where ρ is the system's density matrix, κ_b is the total decay rate of the linear microwave resonator, and γ_1 is the qubit relaxation rate. We do not consider dephasing here since single-shot projections on the z-axis are insensitive to fast frequency fluctuations, and $T_{2\text{echo}}$ in our experiment is much greater than T_1 [56]. The system Hamiltonian derived below is given by,

$$\hat{\mathcal{H}}(t)/\hbar = -\frac{\Delta_q}{2} \hat{\sigma}_z + \omega_b \hat{b}^\dagger \hat{b} + g(t)(\hat{\sigma}_+ \hat{b} + \hat{\sigma}_- \hat{b}^\dagger) + \epsilon_d(t)(\hat{\sigma}_+ e^{-i\Delta_q t} + \hat{\sigma}_- e^{i\Delta_q t}). \quad (3.22)$$

Here, $\Delta_q = \omega_d - \omega_q$ is the drive detuning, $g(t)$ is the time dependent coupling rate, and $\epsilon_d(t)$ is a time dependent resonant qubit drive. The system is initially prepared in its ground state, and the qubit is prepared with a sequence of resonant square wave π -pulses 24 ns in duration and sideband π -swap pulses 91 ns in duration. The resultant time traces are plotted in Fig. 3 in the main text and show excellent agreement with the experimental results.

3.7.5. LONGITUDINAL COUPLING THEORY

QUANTIZATION OF SUPERCONDUCTING QUBIT

As defined in Ref. [37], the Hamiltonian describing an LC oscillator is given by

$$\hat{\mathcal{H}} = 4E_C \hat{n}^2 + \frac{1}{2} E_L \hat{\phi}^2. \quad (3.23)$$

Here, $E_C = e^2/(2C)$, is the charging energy required to add each electron from the Cooper-pair to the island and $E_L = (\Phi_0/2\pi)^2/L$ is the inductive energy, where $\Phi_0 = h/(2e)$ is the magnetic flux quantum. The operator $\hat{n} = \hat{Q}/2e$ is the excess number of Cooper pairs on the island, and $\hat{\phi}$ is the gauge invariant phase across the inductor. Note that these form a pair of canonically conjugate variables $[\hat{\phi}, \hat{n}] = i$.

We can recast this Hamiltonian in the form,

$$\hat{\mathcal{H}} = \hbar\omega_q (\hat{a}^\dagger \hat{a} + 1/2), \quad (3.24)$$

where $\hat{a}(\hat{a}^\dagger)$ is the annihilation (creation) operator of a single excitation of the LC oscillator. The resonance frequency is defined as $\omega_q = \sqrt{8E_C E_L}/\hbar$ in anticipation of this circuit representing the transmon qubit. The charge and flux operators can be written in the form $\hat{n} = i n_{\text{zpf}} (\hat{a} - \hat{a}^\dagger)$ and $\hat{\phi} = \phi_{\text{zpf}} (\hat{a} + \hat{a}^\dagger)$, where $n_{\text{zpf}} = (E_L/(32E_C))^{1/4}$ and $\phi_{\text{zpf}} = (2E_C/E_L)^{1/4}$.

Replacing the inductor with a Josephson junction results in the modified Hamiltonian of the form,

$$\hat{\mathcal{H}} = 4E_C \hat{n}^2 - E_J \cos(\hat{\phi}), \quad (3.25)$$

where $E_C = e^2/(2C_\Sigma)$, $C_\Sigma = C_s + C_j$ is the total capacitance, including the shunt capacitance and the self-capacitance of the Josephson junction, $E_J = I_c \Phi_0/2\pi$ is the Josephson energy, where I_c is the critical current of the junction. In the transmon limit (i.e. $E_J \gg E_C$), the fluctuations of $\hat{\phi}$ are small; therefore, we can expand the potential term, giving

$$-E_J \cos(\hat{\phi}) = \frac{1}{2} E_J \hat{\phi}^2 - \frac{1}{24} E_J \hat{\phi}^4 + O(\hat{\phi}^6). \quad (3.26)$$

The second term results in an anharmonic potential, with an anharmonicity given by $\alpha = \omega_q^{1 \rightarrow 2} - \omega_q^{0 \rightarrow 1}$. For a transmon qubit $\alpha = -E_C$, it is typically 100-300 MHz.

Keeping terms up to fourth order in creation and annihilation operators, the Hamiltonian can be written in the form,

$$\hat{\mathcal{H}}/\hbar = \omega_q \hat{a}^\dagger \hat{a} + \frac{\alpha}{2} \hat{a}^\dagger \hat{a}^\dagger \hat{a} \hat{a}, \quad (3.27)$$

which is the Hamiltonian for a Kerr oscillator. Rather than a single Josephson junction, a pair of junctions are used as a SQUID loop to provide flux tunability to the transmon qubit. The

effective Hamiltonian of such a flux-tunable transmon is given by

$$\hat{\mathcal{H}} = 4E_C \hat{n}^2 - 2E_J |\cos(\varphi_e)| \cos \hat{\phi}. \quad (3.28)$$

Here φ_e is the externally applied magnetic flux. We can define an effective Josephson energy $\tilde{E}_J(\varphi_e) = 2E_J |\cos(\varphi_e)|$.

LONGITUDINAL HAMILTONIAN

Considering the linear resonator and the coherent drive tones, the total Hamiltonian of the device is

$$\hat{\mathcal{H}} = 4E_C \hat{n}^2 - \tilde{E}_J(\varphi_e) \cos \hat{\phi} + \hbar \omega_b \hat{b}^\dagger \hat{b} + \hat{\mathcal{H}}_d. \quad (3.29)$$

Here, $\hat{b}$ ($\hat{b}^\dagger$) are the annihilation (creation) operators for the linear resonator. During operation, the external flux bias is fixed at a single value, and due to the mutual inductive coupling between the linear circuit and the SQUID loop, the external flux is modulated by the linear resonator. We can, therefore, write the Josephson energy in the form

$$\tilde{E}_J(\varphi_e) = \bar{E}_J + \frac{d\bar{E}_J}{d\varphi_e} \frac{d\varphi_e}{dI} I_{zpf} (\hat{b} + \hat{b}^\dagger). \quad (3.30)$$

Here, $\bar{E}_J$ is the Josephson energy at the flux-bias point, $I_{zpf} = \sqrt{\hbar \omega_b / 2L_b}$, where L_b is the inductance of the linear resonator. We can simplify this equation by defining a coupling constant in the form

$$\hbar \tilde{g}_0 = \frac{d\bar{E}_J}{d\varphi_e} \frac{d\varphi_e}{dI} I_{zpf}. \quad (3.31)$$

Therefore, the Hamiltonian can be written as,

$$\hat{\mathcal{H}} = 4E_C \hat{n}^2 - \bar{E}_J \cos \hat{\phi} - \hbar \tilde{g}_0 \cos \hat{\phi} (\hat{b} + \hat{b}^\dagger) + \hbar \omega_b \hat{b}^\dagger \hat{b} + \hat{\mathcal{H}}_d. \quad (3.32)$$

At this stage, we can expand the cosine of the gauge invariant flux using Eq. 3.26, giving

$$\hat{\mathcal{H}} = 4E_C \hat{n}^2 + \bar{E}_J \left(\frac{1}{2} \hat{\phi}^2 - \frac{1}{24} \hat{\phi}^4 \right) + \hbar \tilde{g}_0 \frac{1}{2} \hat{\phi}^2 (\hat{b} + \hat{b}^\dagger) + \hbar \omega_b \hat{b}^\dagger \hat{b} + \hat{\mathcal{H}}_d. \quad (3.33)$$

Keeping only first-order terms in the interaction and using the definitions $\hat{\phi} = \phi_{zpf} (\hat{a} + \hat{a}^\dagger)$ and $\phi_{zpf} = (2E_C / \bar{E}_J)^{1/4}$ we arrive at the Hamiltonian based on creation and annihilation operators. The Hamiltonian can be written in the form,

$$\hat{\mathcal{H}} / \hbar = \omega_q \hat{a}^\dagger \hat{a} + \frac{\alpha}{2} \hat{a}^\dagger \hat{a}^\dagger \hat{a} \hat{a} + \omega_b \hat{b}^\dagger \hat{b} + \frac{\tilde{g}_0}{2} \left(\frac{2E_C}{\bar{E}_J} \right)^{1/2} (\hat{a} + \hat{a}^\dagger)^2 (\hat{b} + \hat{b}^\dagger) + \hat{\mathcal{H}}_d / \hbar. \quad (3.34)$$

Where, we will define $g_0 = \tilde{g}_0 (2E_C / \bar{E}_J)^{1/2}$. Moreover, the drive Hamiltonian can be written in the form,

$$\hat{\mathcal{H}}_d = \hbar \epsilon_d (\hat{a} e^{i\omega_d t} + \hat{a}^\dagger e^{-i\omega_d t}) \quad (3.35)$$

Here, $\epsilon_d = \sqrt{\mathcal{P} \kappa_e / \hbar \omega_d}$ is the coherent drive strength. We can now transform the Hamiltonian into a frame rotating at the drive frequency using the unitary transformation $\hat{\mathcal{U}} = e^{i\omega_d \hat{a}^\dagger \hat{a} t}$,

where $\hat{\mathcal{H}}_{\text{rot}} = \hat{\mathcal{U}}\hat{\mathcal{H}}\hat{\mathcal{U}}^\dagger + i\hbar(d\hat{\mathcal{U}}/dt)\hat{\mathcal{U}}^\dagger$, keeping only co-rotating terms we can express the new Hamiltonian as

$$\hat{\mathcal{H}}/\hbar = -\Delta_q \hat{a}^\dagger \hat{a} + \frac{\alpha}{2} \hat{a}^\dagger \hat{a}^\dagger \hat{a} \hat{a} + \omega_b \hat{b}^\dagger \hat{b} + g_0 \hat{a}^\dagger \hat{a} (\hat{b} + \hat{b}^\dagger) + \epsilon_d (\hat{a} + \hat{a}^\dagger). \quad (3.36)$$

Here $\Delta_q = \omega_d - \omega_q$. We can first study the steady-state classical dynamics of the Kerr oscillator. The classical equation of motion is given by

$$\frac{d}{dt} \langle \hat{a} \rangle = (i\Delta_q - \gamma_q/2) \langle \hat{a} \rangle - i\alpha |\langle \hat{a} \rangle|^2 \langle \hat{a} \rangle - ig_0 (\langle \hat{b} \rangle + \langle \hat{b} \rangle^*) \langle \hat{a} \rangle - \epsilon_d. \quad (3.37)$$

The coupling to the linear resonator mode has the effect of modifying the Kerr nonlinearity as described in [57], with an effective Kerr nonlinearity of

$$\alpha_{\text{eff}} = \alpha - \frac{2g_0^2 \omega_b}{\omega_b^2 + \gamma_b^2/4}. \quad (3.38)$$

Where γ_b is the linewidth of the linear mode. Using the parameters from our device, the correction to the Kerr nonlinearity is much less than the Kerr nonlinearity. Therefore, the equation of motion can be written in the form,

$$\frac{d}{dt} \langle \hat{a} \rangle = (i\Delta_q - \gamma_q/2) \langle \hat{a} \rangle - i\alpha |\langle \hat{a} \rangle|^2 \langle \hat{a} \rangle - \epsilon_d. \quad (3.39)$$

In steady-state (i.e. $\dot{\langle \hat{a} \rangle} = 0$), the cubic equation has one real root, but for a large enough input drive, a bistable regime exists. The cubic equation has the form,

$$\bar{n}_q (-\Delta_q + \alpha \bar{n}_q)^2 - \epsilon_d^2 = 0. \quad (3.40)$$

Here, we have defined $\bar{n}_q = |\langle \hat{a} \rangle|^2$ the steady-state amplitude of the Kerr oscillator.

PHOTON NUMBER CALIBRATION

To calibrate the number of photons in the qubit due to the sideband drive tone, we use the AC-Stark shift. The driving term applied on the red-sideband results in an AC-Stark shift of the Kerr oscillator. Since the drive tone detuning is large ($\Delta_q = \omega_d - \omega_q \approx \omega_b = 4.35$ GHz), we must consider the Bloch-Siegert shift due to the counter-rotating terms, which is described in Ref. [33]. The AC-Stark shift considering the co- and counter-rotating terms is given by,

$$\delta\omega_q = \frac{1}{2} \epsilon_d^2 \alpha \left(\frac{1}{\Delta_q^2} + \frac{2}{|\Delta_q| \Sigma} + \frac{1}{\Sigma^2} \right). \quad (3.41)$$

Where $\Sigma = \omega_q + \omega_d$, α is the transmon anharmonicity and $\epsilon_d = \sqrt{\mathcal{P} \kappa_e / \hbar \omega_d}$. We can re-arrange this equation by writing it in the form,

$$\epsilon_d^2 = \frac{2\delta\omega_q}{\alpha} \left(\frac{1}{\Delta_q^2} + \frac{2}{|\Delta_q| \Sigma} + \frac{1}{\Sigma^2} \right)^{-1}. \quad (3.42)$$

Therefore, we can calibrate the average photon number in the Kerr oscillator using the AC-Stark shift. Combining Eq. 3.40 and Eq. 3.42, the average photon number can be found by solving the cubic equation,

$$\bar{n}_q(-\Delta_q + \alpha \bar{n}_q)^2 - \frac{2\delta\omega_q}{\alpha} \left(\frac{1}{\Delta_q^2} + \frac{2}{|\Delta_q|\Sigma} + \frac{1}{\Sigma^2} \right)^{-1} = 0. \quad (3.43)$$

The first real solution to this equation is the photon population of the Kerr-oscillator, $\langle \hat{a}^\dagger \hat{a} \rangle = \bar{n}_q$. Note that treating the qubit as a two-level system at this step underestimates the sideband population and, therefore, will overestimate a g_0 .

3

3.7.6. LINEARIZED COUPLING

At this point, to continue our analysis, we aim to linearize the theory about the sideband drive applied to the qubit. From the AC-Stark shift calibration above, we found that for all drive powers, the average occupancy $|\langle \hat{a} \rangle|^2$ is much less than a single photon on average. Therefore, we will perform our analysis assuming a harmonic oscillator, making the two-level approximation once we have arrived at a linearized theory.

SIDEBAND LINEARIZATION

If we begin with the Hamiltonian given by Eq. 3.34, rewriting it in the form,

$$\hat{\mathcal{H}}/\hbar = \omega_q \hat{a}^\dagger \hat{a} + \frac{\alpha}{2} \hat{a}^\dagger \hat{a}^\dagger \hat{a} \hat{a} + \omega_b \hat{b}^\dagger \hat{b} + \frac{g_0}{2} (\hat{a}^2 + (\hat{a}^\dagger)^2 + 2\hat{a}^\dagger \hat{a} + 1)(\hat{b} + \hat{b}^\dagger) + (\epsilon_d^* \hat{a} e^{i\omega_d t} + \epsilon_d \hat{a}^\dagger e^{-i\omega_d t}). \quad (3.44)$$

In the experimentally relevant situation where the sideband drive is always weak such that $|\langle \hat{a} \rangle|^2 \ll 1$, we will drop the Kerr term as the Kerr oscillator is well approximated by a harmonic oscillator in this weak driving regime. This results in the Hamiltonian

$$\hat{\mathcal{H}}/\hbar = \omega_q \hat{a}^\dagger \hat{a} + \omega_b \hat{b}^\dagger \hat{b} + g_0 \hat{a}^\dagger \hat{a} (\hat{b} + \hat{b}^\dagger) + (\epsilon_d^* \hat{a} e^{i\omega_d t} + \epsilon_d \hat{a}^\dagger e^{-i\omega_d t}). \quad (3.45)$$

We can simultaneously transform into a frame rotating at the drive frequency and displaced by the steady-state amplitude $\langle \hat{a} \rangle$. This is equivalent to applying the unitary

$$\hat{\mathcal{D}}\hat{\mathcal{U}} = e^{\langle \hat{a} \rangle^* \hat{a} - \langle \hat{a} \rangle \hat{a}^\dagger} e^{i\omega_d \hat{a}^\dagger \hat{a} t}. \quad (3.46)$$

The resulting Hamiltonian in the transformed frame is

$$\hat{\mathcal{H}}/\hbar = -\Delta_q \hat{a}^\dagger \hat{a} + \omega_b \hat{b}^\dagger \hat{b} + g_0 (\langle \hat{a} \rangle^* \hat{a} + \langle \hat{a} \rangle \hat{a}^\dagger) (\hat{b} + \hat{b}^\dagger) + g_0 \hat{a}^\dagger \hat{a} (\hat{b} + \hat{b}^\dagger). \quad (3.47)$$

This transformation has separated the interaction into its linear and non-linear components. In this way, we have *only* performed unitary transformations resulting in static shifts of the qubit's equilibrium frequency. Moreover, the last term in this Hamiltonian results in a non-linear two-photon interaction, as described in Ref. [36]. This term can be dropped since it does not affect the dynamics in the frame of our sideband drive. However, this term is crucial for generating Schrodinger cat states via the non-linearity of the longitudinal interaction. We can write our Hamiltonian in the simplified form,

$$\hat{\mathcal{H}}/\hbar = -\Delta_q \hat{a}^\dagger \hat{a} + \omega_b \hat{b}^\dagger \hat{b} + g(\hat{a} + \hat{a}^\dagger)(\hat{b} + \hat{b}^\dagger). \quad (3.48)$$

We define the enhanced coupling rate as $g = |\langle \hat{a} \rangle| g_0$ and have assumed it is real. Performing the RWA and approximating the harmonic oscillator as a two-level system, we can write the final Hamiltonian as,

$$\hat{\mathcal{H}}/\hbar = -\frac{\Delta_q}{2} \hat{\sigma}_z + \omega_b \hat{b}^\dagger \hat{b} + g(\hat{\sigma}_+ \hat{b} + \hat{\sigma}_- \hat{b}^\dagger). \quad (3.49)$$

We arrive at the drive-induced Jaynes-Cummings Hamiltonian. Following standard quantum optics techniques [58], one can derive the N excitation Rabi frequency given by $2g_N = 2\sqrt{N}g$.

NORMAL-MODE SPECTRUM

Considering the Hamiltonian given in Eq. 3.49, we can derive the eigenfrequencies to fit the normal-mode spectra. If we consider the quantity $\hat{N} = \hat{a}^\dagger \hat{a} + \hat{\sigma}_+ \hat{\sigma}_-$ which defines the total number of excitations within the system, in the absence of decay. We notice that $[\hat{\mathcal{H}}, \hat{N}] = 0$, thus is a conserved quantity of the system. Therefore, we can consider the block-diagonal element of the Hamiltonian defined by the total number of excitations spanning the subspace $\{|g, n\rangle, |e, n-1\rangle\}$. This block diagonal element can be written in the form,

$$\hat{\mathcal{H}}/\hbar = \begin{pmatrix} -\Delta_{qb} + n\omega_b & g\sqrt{n} \\ g\sqrt{n} & n\omega_b \end{pmatrix}. \quad (3.50)$$

Where $\Delta_{qb} = \omega_b + \Delta_q$, recalling that in the rotating frame $\Delta_q = \omega_d - \omega_q \approx -\omega_b$; therefore, this is the qubit-linear resonator detuning in the rotating frame. Following some algebra, it can be shown that the energy eigenvalues of this block diagonal Hamiltonian have the form [59],

$$E_{n,\pm} = n\hbar\omega_b - \hbar\frac{\Delta_{qb}}{2} \pm \frac{\hbar}{2}\sqrt{\Delta_{qb}^2 + 4ng^2}, \quad (3.51)$$

from which we can recover the n -quanta Rabi frequency $2g_n = 2\sqrt{n}g$. If we consider the lowest energy manifold of the Jaynes-Cummings ladder, we can determine the expected spectroscopic eigenfrequencies. For $n = 1$, we find the frequency spectrum given by,

$$\omega_{\pm} = \omega_b - \frac{\Delta_{qb}}{2} \pm \frac{1}{2}\sqrt{\Delta_{qb}^2 + 4g^2}. \quad (3.52)$$

This can be written in the more enlightening form,

$$\omega_{\pm} = \frac{\omega_b - \Delta_q}{2} \pm \frac{1}{2}\sqrt{\Delta_{qb}^2 + 4g^2}. \quad (3.53)$$

Using Eq. 3.53, we can fit the experimental data to extract g as a function of sideband drive power, as shown in the main text Fig. 2. Using the extracted value of g and the photon number extracted from Eq. 3.43, we can determine the single-photon coupling rate g_0 , as shown in the main text. It is important to note that this solution *only* holds when the system is within the lowest energy manifold of the Jaynes-Cummings ladder.

Strong Intrinsic Longitudinal Coupling in Circuit Quantum Electrodynamics

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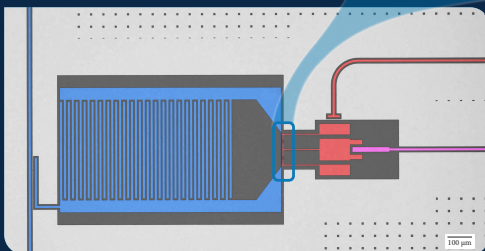
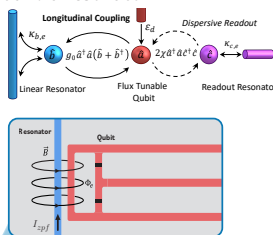
Longitudinal coupling between a qubit and a resonator

- A parametric light-matter interaction, also known as longitudinal coupling in cQED, between a superconducting qubit and resonator is realized.
- Zero-point current fluctuations through the resonator generate a flux through the SQUID of the qubit.
- Frequency of the qubit is parametrically modulated by the generated flux.
- Coupling Hamiltonian is given by

$$\hat{H}_{\text{coupling}} = \hbar g_0 \hat{a}^\dagger \hat{a} (\hat{b}^\dagger + \hat{b}),$$

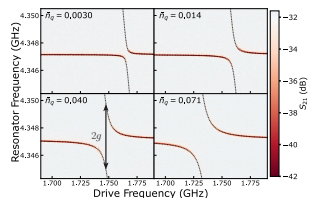
with

$$g_0 = \frac{\partial \omega_q}{\partial \Phi_e} \Phi_{\text{ZPF}}.$$



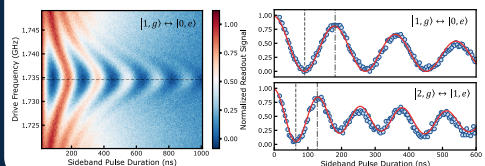
Drive-Induced Jaynes-Cummings

- Drive-induced Jaynes-Cummings interaction by applying a strong sideband drive.
- The coupling rate is $g = g_0 \sqrt{n_q}$ with n_q the average sideband steady-state population of the qubit.
- We observe a series of avoided level crossings in the resonator as we sweep the sideband drive frequency.



Vacuum Rabi Oscillations

- Coherence of interaction is demonstrated by measuring vacuum Rabi oscillations.
- The qubit is excited and a sideband pulse of variable duration is applied to drive the coherent exchange of an excitation between qubit and resonator.
- Drive-induced Jaynes-Cummings ladder is confirmed by preparing the resonator in its first excited state and subsequently measuring Rabi oscillations.
- Exchange rate increased by a factor of $\sqrt{2}$ as predicted by Jaynes-Cummings model.



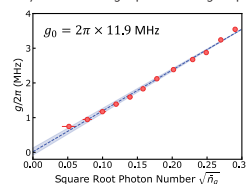
Outlook

In this work we demonstrated a strong intrinsic longitudinal coupling between a superconducting qubit and resonator. This novel interaction will enable several promising applications in the future:

- Interfacing cQED hardware with cavity optomechanical systems for testing theories of gravitational decoherence of quantum objects.
- Testing theories of single-photon strong coupling, such as dissipative cat state generation.
- A new platform for superconducting quantum information processing with a potential for high connectivity and no residual interactions.

Strong Intrinsic Longitudinal Coupling

- Coupling rate g measured for a range of sideband drive powers.
- Qubit experiences increasing AC-Stark shift for increasing drive power, providing an absolute calibration of n_q .
- Single-photon coupling rate is extracted from slope of g as a function of $\sqrt{n_q}$.
- We extract $g_0 = 2\pi \times 11.9$ MHz, larger than the decay rates, placing our system in the single-photon strong coupling regime.



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Figure 3.9: Poster for the paper: *Strong Intrinsic Longitudinal Coupling in Circuit Quantum Electrodynamics*, *Physical Review Letters*, 134, 153603, (2025). Presented at NWO Physics 2025 and Helgoland 2025.

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4

FABRICATION AND CHARACTERIZATION OF A SNAIL PARAMETRIC AMPLIFIER

Parametric amplifiers are widely used in superconducting quantum experiments to enable low-noise microwave signal detection. In this work, we present the design, fabrication, and characterization of flux-tunable parametric amplifiers based on the Superconducting Nonlinear Asymmetric Inductive eLement (SNAIL). The devices consist of a resonator-based amplifier that enables three-wave mixing while suppressing unwanted Kerr nonlinearities. Using wafer-scale fabricated devices, we characterize the linear and nonlinear response of the amplifiers at millikelvin temperatures. We demonstrate stable, high-gain parametric amplification over a wide flux-tunable frequency range, investigate the gain-bandwidth trade-off, and characterize the dynamic range through gain-compression measurements. The results provide an experimental characterization of SNAIL-based parametric amplifiers across their tunable operating range. Together, these measurements characterize the operating regime and performance limits of the SNAIL-based parametric amplifiers realized in this work.

4.1. INTRODUCTION

Parametric amplifiers have become an essential component in modern quantum microwave experiments. Their ability to provide large gain with added noise near the quantum limit, makes them well suited for detecting weak microwave signals near the single photon limit. They are used in high-fidelity qubit readout [1–7], in microwave optomechanics [8, 9], for generating squeezed states [10–15], and in the search for axion dark matter [16–21]. Parametric amplification was first introduced to the microwave domain in the 1960s, where early implementations were realized using variable-capacitance (varactor) diodes [22]. These early devices established the principle of parametric gain for microwave signals, but their performance was limited by the dissipation and noise of the used semiconductor components. The rise of superconducting circuits briefly revived the interest in parametric amplification in the 1980s, which resulted in the first Josephson junction based parametric amplifier demonstrated by Yurke *et al.* [23]. This was the first demonstration of parametric gain and quantum squeezing in a Josephson parametric amplifier (JPA). However, they did not become widely adopted because of their small frequency tunability on the order of 1 MHz. Furthermore, at that time, there were also not yet many applications that required near-quantum-limited noise performance.

This changed with the rise of superconducting quantum information processing, where dispersive qubit readout brought the need for high-gain and low noise microwave amplification [24]. Conventionally used high-electron-mobility transistor (HEMT) amplifiers typically added 10–20 photons of noise per signal photon at 5 GHz [25], which limited the fidelity of qubit readout. This interest from the quantum community resulted in the first practical widely tunable JPA being developed [26]. Today, a wide variety of JPAs exist, current pumped JPAs (four-wave mixing) [26–28] and flux pumped JPAs (three-wave mixing), which offered improved isolation by operating the signal and pump tone at different frequencies [29, 30]. JPAs capable of operating in large magnetic field [31, 32] and at frequencies above 20 GHz [33] are now also available. Modern impedance-engineered amplifier designs can be complex to engineer, but also offer large bandwidths [6, 34–36].

In parallel with these developments, the simpler resonator-based parametric amplifier architectures have remained highly relevant. A major advancement was the introduction of the Superconducting Nonlinear Asymmetric Inductive eLement (SNAIL) [37–39]. The SNAIL provides a tunable, asymmetric inductive nonlinearity that enables three-wave mixing while simultaneously suppressing unwanted Kerr (fourth-order) nonlinearities. This suppression significantly improved the dynamic range that limited SQUID-based JPAs by minimizing pump-induced resonator shifts. Here, we report the wafer-scale fabrication and characterization of SNAIL-based parametric amplifiers (SPAs) developed in our lab. Our aim is to realize devices with a flux-tunable bandwidth exceeding 1 GHz, using multiple SPA designs that collectively span the 4–8 GHz frequency range. We present the gain and bandwidth performance across the full flux-tunable operating window, and we highlight the achieved dynamic range. These devices are intended for applications in high-fidelity qubit readout and in microwave optomechanical experiments.

4.2. DEVICE DESCRIPTION

The SPA consists of a $\lambda/2$ microstrip resonator that is capacitively coupled to a feedline on both ends, forming the signal and pump ports. Fig. 4.1 (a) shows an overview of the complete chip. The central section of the resonator contains an array of 20 SNAILS, as shown in the zoom in of Fig. 4.1 (b). Each SNAIL is made up of a superconducting loop with three large Josephson junctions and one smaller junction, providing a junction asymmetry ratio α . The Al/AIO_x/Al Josephson junctions were fabricated using the Dolan bridge double-angle evaporation technique [40]. A detailed view of a single SNAIL is shown in Fig. 4.1 (c), where the junction areas are highlighted in red. The device is fabricated on a 325 μm thick high-resistivity Si substrate. The backside of the chip is coated with a 200 nm gold layer, which serves as the non-superconducting ground plane for the microstrip geometry.

The microstrip resonator was designed to support an amplification bandwidth of approximately 20 MHz at 20 dB gain, which corresponds to a bare resonator linewidth on the order of 200 MHz. In addition, it is necessary that internal dissipation remains small compared to the external coupling rate, so that the amplifier performance is not limited by internal loss. To assess the intrinsic losses of the microstrip resonator, we fabricated a test resonator with a geometry identical to that of the SPA resonator, but without the signal line and with the SNAIL chain replaced by a short. The measured total linewidth of this test resonator was $\kappa_{\text{tot}}/2\pi = 14.4$ MHz. Because the resonator is overcoupled, a precise separation of the external and internal decay rates is not highly accurate [41]. Nevertheless, fitting the response yields values of $\kappa_{\text{ext}}/2\pi \approx 10$ MHz and $\kappa_{\text{int}}/2\pi \approx 4$ MHz, indicating that internal losses are small. A more precise quantification of the internal loss rate would require further dedicated measurements, but these results indicate that internal dissipation is not expected to limit the amplifier performance.

Fig. 4.1 (d) shows the chip mounted inside a gold-plated copper sample holder. The gold backplane of the chip is glued to the copper PCB using silver paste. The pump and signal port are wire-bonded to two microwave launchers on the PCB that connect to SMA ports on the package. An external coil magnet is mounted on the back side of the copper box to apply a uniform magnetic flux Φ to the SNAILS. The entire assembly is enclosed in a compact cryoperm shield. The microwave package and magnetic shielding are based on a compact design introduced by the Yale group [39].

The SPA was characterized mounted at the MXC plate of a Bluefors dilution refrigerator with a base temperature of approximately 25 mK. Fig. 4.1 (e) illustrates a simplified schematic of how the SPA is integrated into the measurement setup. The complete wiring diagram can be found in Supplementary Information Fig. 4.9. The input signal is routed through a circulator into the input port of the SPA. The amplified signal is reflected back to the circulator, which then routes the signal to the rest of the output amplification chain. The pump signal is injected through the separate pump port on the opposite side of the SPA packaging. Further details on the fabrication process and the cryogenic measurement setup are provided in the Supplementary Information.

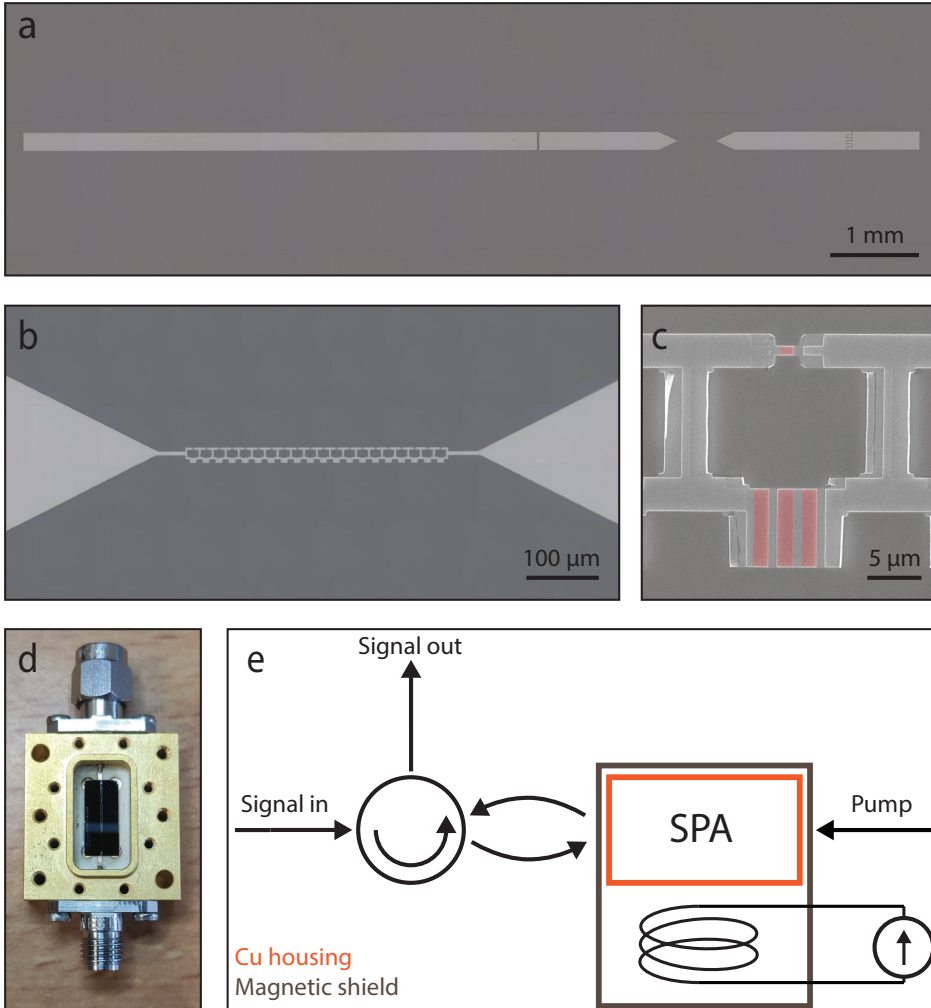


Figure 4.1: Device layout and packaging. (a) Optical image of the full SPA chip. The $\lambda/2$ microstrip resonator is capacitively coupled to two independent feedlines, which serve as the signal and pump ports. (b) Magnified view of the central section of the resonator containing an array of 20 SNAILs. (c) SEM image of a single SNAIL. Each SNAIL consists of a superconducting loop containing three large Josephson junctions and one smaller junction as highlighted in red. (d) Photograph of the SPA chip mounted in the gold-plated copper package. The chip is wire-bonded to microwave launchers on the PCB that are connected to SMA ports to interface with the cryogenic wiring. (e) A simplified schematic of the cryogenic measurement setup.

4.3. EXPERIMENTAL RESULTS

FREQUENCY TUNABLE NONLINEAR RESONATOR

To characterize the SPA, we first examine its linear response by measuring the reflection from the signal port. The circulator directs the reflected signal to a separate output line, such that the complex transmission coefficient S_{21} between input and output ports can be measured with a vector network analyzer (VNA). Fig. 4.2 (a,b) show the measured phase and magnitude of the reflected signal as a function of probe frequency and applied magnetic flux.

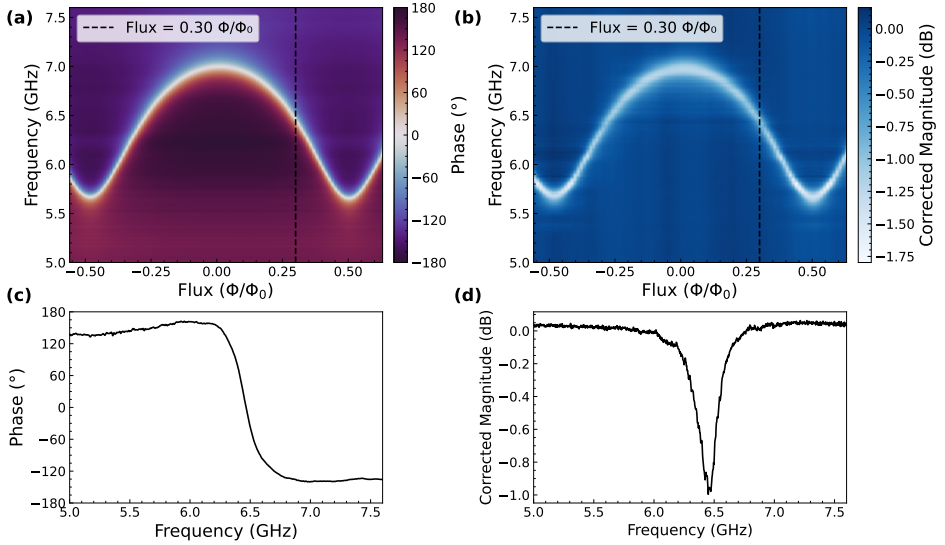


Figure 4.2: Linear response of the SNAIL parametric amplifier as a function of flux. (a) Measured phase response of the SPA as a function of the applied magnetic flux and probe frequency. The resonance frequency shows a periodic modulation with flux. The dashed line indicates the flux bias point at $\Phi/\Phi_0 = 0.30$ where the line cut shown in panel (c) is taken. (b) Corresponding magnitude response after background correction, which removed the contribution of the slowly varying background. (c) Line cut of the phase response at a flux bias of $\Phi/\Phi_0 = 0.30$, showing that the phase of the reflected signal shifts by approximately 180° as the probe frequency passes through the resonance frequency. (d) Line cut of the corrected magnitude response at the same flux bias point.

The magnetic flux through the SNAIL array is tuned by sending a current through the external superconducting coil. Before plotting the magnitude data, a smooth background was subtracted to remove slow frequency dependent variations in the cable and component transmission. The background was obtained by averaging the off-resonant response over all flux values, excluding a frequency window around the resonance at each flux point. The full background-subtraction procedure is described in supplementary information 5.5.4. To interpret the residual amplitude response, we can introduce the coupling efficiency η of the unpumped resonator, that is described by

$$\eta = \frac{\kappa_{\text{ext}}}{\kappa_{\text{tot}}} = \frac{\kappa_{\text{ext}}}{\kappa_{\text{ext}} + \kappa_{\text{int}}}, \quad (4.1)$$

where κ_{int} and κ_{ext} are the internal and external coupling rates. In the ideal limit of a lossless perfectly overcoupled resonator where $\eta = \kappa_{\text{ext}}/\kappa_{\text{tot}} = 1$, the on-resonance magnitude would remain unity as the signal is fully reflected. The presence of a small amplitude dip in the measured amplitude response shows that some energy is lost inside the resonator, indicating that $\eta < 1$ and that non-negligible internal loss is present in the device. Panels 4.2 (c,d) show line cuts of the phase and corrected magnitude responses at a flux bias of $\Phi/\Phi_0 = 0.30$. The phase of the reflected signal shifts by approximately 180° , characteristic of a resonant reflection in the overcoupled limit ($\eta > 0.5$), while the corresponding magnitude trace exhibits a small dip at the same frequency.

To quantify the resonance frequency and loss rates of the SPA as functions of applied flux, we fit the measured complex data at each bias point using a single-port reflection model. Fig. 4.3 (a) shows the resonance frequency extracted from the fit as a function of magnetic flux Φ/Φ_0 .

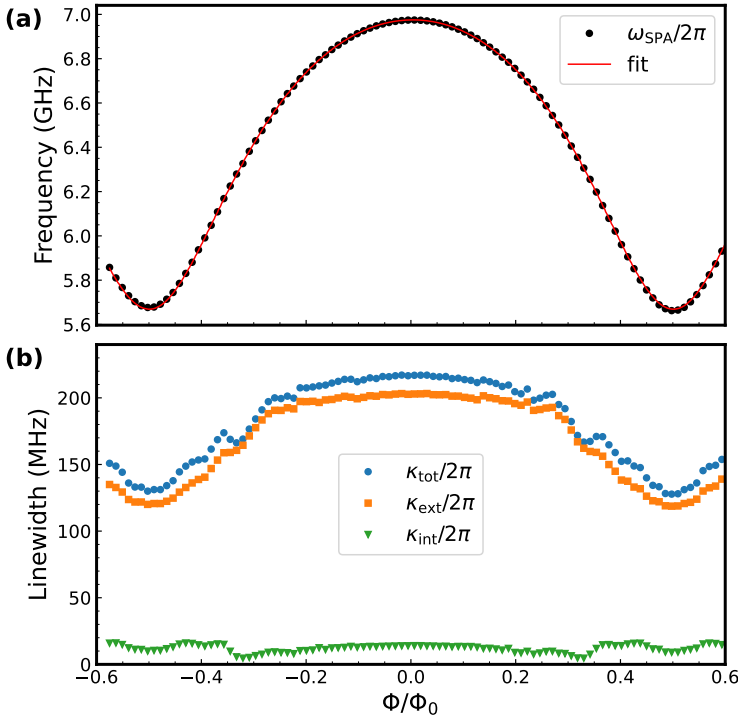


Figure 4.3: Extracted resonance frequency and linewidth as a function of flux. (a) The resonance frequency ω_{SPA} obtained from a fit to the complex reflection response as a function of the applied magnetic flux Φ/Φ_0 . The SPA exhibits a periodic modulation with flux and is tunable from 5.66 GHz to 6.97 GHz, corresponding to a tunable range of approximately 1.3 GHz. (b) Corresponding linewidth extracted from the same fit as a function of flux.

The resonance frequency varies periodically with flux, consistent with the expected modulation of the SNAIL inductance. The SPA is tunable from 5.66 GHz to 6.97 GHz, corresponding to a tunable range of approximately 1.3 GHz. We also fit the extracted resonance frequencies to the theoretical SPA model introduced in [38]. To fit our data to this model, we fix the number of SNAILs to $M = 20$, and set the impedance of the microstrip geometry to the designed value of $Z_a = 50 \Omega$. From this we extract the device parameters, and find the inductance of a single Large junction $L_J = 26.6$ pH, $\alpha = 0.10$ and $\omega_0/2\pi = 10.1$ GHz, which is the fundamental frequency of a square microstrip resonator in the absence of the SNAIL array. A detailed description of the fitting procedure is provided in Supplementary Information 4.5.6.

Fig. 4.3 (b) shows the extracted linewidths at each flux point. The fits indicate that the SPA is overcoupled, with a coupling efficiency $\eta = \kappa_{ext}/\kappa_{tot} \approx 0.9$. The internal linewidth κ_{int} is on the order of 10-20 MHz, while the total linewidth is dominated by the external coupling. Small fluctuations in the extracted κ_{int} are observed, which we suspect to arise primarily due to systematic uncertainties in the fit rather than intrinsic physical variations. These are likely caused by impedance variations in the measurement chain. In particular, interference between different paths due to cable reflections, to which the reflection geometry is particularly sensitive. This can lead to an asymmetric Fano lineshape in the resonator response, which distorts the extracted internal and external coupling rates [41]. The total linewidth κ_{tot} is a more reliable quantity to interpret. We observe that the linewidth changes as a function of flux and reaches a maximum value on the order of 200 MHz around zero flux. This is similar to other reported SPA devices [39]. The small flux-dependent variations in κ_{tot} on top of the background, for example the feature around $\Phi/\Phi_0 \approx 0.32$, are attributed to interactions of the SPA resonance with impedance features in the measurement chain, such as frequency dependent variation in the environmental impedance of the setup arising from cable resonances. This can also be observed in the remaining distortions of the line shape in the corrected magnitude linecut of Fig. 4.2 (d).

AMPLIFIER CHARACTERIZATION.

Having established the linear response properties of the SPA resonator, we now characterize the performance as a three-wave mixing parametric amplifier. A strong microwave pump tone is applied through the dedicated pump port, with the pump frequency set to $\omega_p = 2\omega_{SPA}$. This drives the system into the three-wave-mixing regime, in which the signal and idler sidebands both appear at the resonator frequency ω_{SPA} [14].

To quantify the amplifier gain, we measure the complex transmission coefficient S_{21} between the signal input and output ports while varying the pump power at a fixed flux bias point. We note the distinction between power gain and linear gain. Power gain refers to the ratio of output to input power, and is often expressed in decibels. Linear gain, by contrast, describes the ratio of output to input signal amplitude. While both quantify amplification, power gain is more commonly used in the context of microwave amplifiers, as it directly corresponds to the quantities measured in the experiment. The power gain G can be defined as

$$G = \frac{P_{out}}{P_{in}} = \left| \frac{S_{21}^{pump\ on}}{S_{21}^{pump\ off}} \right|^2, \quad (4.2)$$

where P_{in} and P_{out} are the input and output powers. The ratio compares the transmitted power when the pump is on to that when it is off. In this way, the unpumped response is also used as a reference to normalize the frequency-dependent background of the measurement chain. The power gain in dB is then given by

$$G_{dB} = 10 \log_{10} \left(\frac{P_{out}}{P_{in}} \right). \quad (4.3)$$

Fig. 4.4 illustrates how the amplifier properties of the SPA scale at a fixed flux bias point when

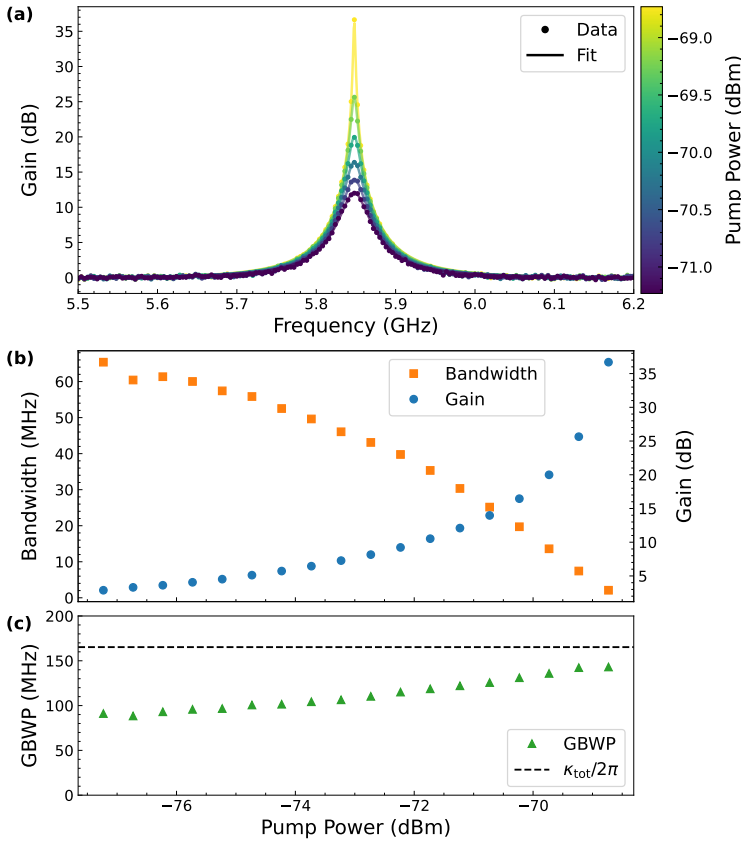


Figure 4.4: Gain, bandwidth, and gain-bandwidth product as a function of pump power at a fixed flux point. (a) Measured gain of the SPA for several pump powers. The resonance frequency is fixed at $\omega_{SPA}/2\pi = 5.847$ GHz, and the pump is applied at $\omega_p = 2\omega_{SPA}$. (b) Gain (blue circles) and the corresponding 3 dB bandwidth (orange squares) extracted from Lorentzian fits to the gain curves. As the pump power increases, the gain grows while the bandwidth decreases, illustrating the gain-bandwidth trade-off. (c) The gain-bandwidth product (GBWP, green triangles) obtained from the fitted gain and bandwidth values. The dashed line shows the previously measured resonator linewidth $\kappa_{tot}/2\pi = 165$ MHz at this flux point.

the pump power is increased. Fig. 4.4 (a) shows the measured gain at a frequency of $\omega_{\text{SPA}}/2\pi = 5.847$ GHz with the pump applied at $\omega_p = 2\omega_{\text{SPA}}$. A series of different pump powers are plotted, where the pump power at the device is calibrated using the method described in supplementary information 4.5.5. It can be observed that when the pump power increases, the gain also increases while the bandwidth becomes narrower. To further quantify these changes, we fit a Lorentzian to each individual gain trace. From these fits, we obtain the maximum gain and the corresponding 3 dB bandwidth, which is defined as the full width at half maximum (FWHM). Fig. 4.4 (b) shows the small-signal gain and bandwidth as a function of pump power. As the pump power increases, the gain rises while the bandwidth decreases, illustrating the gain–bandwidth trade-off. At this operating point, we obtain 20 dB of gain with a bandwidth of 13.6 MHz, which is similar to previously reported values for SPA's [38, 39]. We note that only the stable operating regime is plotted, at higher pump powers the SPA becomes unstable and the gain rapidly collapses. The combined effect of these opposing trends is captured in Fig. 4.4 (c), where we plot the gain–bandwidth product (GBWP) that can be defined as

$$\text{GBWP} = \sqrt{G} \cdot \text{BW}, \quad (4.4)$$

where G is the measured power gain close to $G_{\text{dB}} = 20$, and BW is the corresponding 3 dB bandwidth. We observe that the GBWP gradually increases as a function of pump power, since the gain grows faster than the bandwidth decreases. For comparison, we also indicate the resonator linewidth of $\kappa_{\text{tot}}/2\pi = 165$ MHz at this flux point. It can be seen that the GBWP approaches, but remains slightly below $\kappa_{\text{tot}}/2\pi$. The GBWP is often expected to be approximately constant and limited by the total resonator linewidth, such that $\text{GBWP} \approx \kappa_{\text{tot}}/2\pi$ [27, 42]. However, our data indicates a clear dependence on pump power. Recent studies have shown that the GBWP does not necessarily remain constant across all pump powers, which could explain our observations [43, 44]. In principle, increasing the resonator linewidth $\kappa_{\text{tot}}/2\pi$ can increase the achievable GBWP. However, reaching the same gain with a larger linewidth requires a higher pump amplitude, which ultimately limits the maximum gain of the amplifier.

Building on this understanding at a fixed flux point, we next examine how the SPA performance varies over its tunable frequency range. Fig. 4.5 (a) shows the measured gain as a function of frequency for a series of flux bias points spanning $\Phi/\Phi_0 \approx 0.16$ to $\Phi/\Phi_0 \approx 0.46$. Within this range, the SPA produces at least 20 dB of gain from approximately 5.7 GHz to 6.8 GHz, corresponding to a flux-tunable operational bandwidth of around 1.1 GHz. This tuning range is roughly two orders of magnitude larger than the instantaneous bandwidth. Outside this frequency range, the gain drops below 20 dB as the three-wave mixing nonlinearity g_3 is reduced. Fig. 4.5 (b) shows the extracted 3 dB bandwidth, GBWP, and resonator linewidth as a function of the SPA frequency. Across the tunable range, we obtain instantaneous bandwidths between 7 and 27 MHz, from which the GBWP is calculated. These values are taken from the first pump-power that provides at least 20 dB of gain at each flux point. We note that for this data set, the pump power steps turned out to be relatively coarse at 0.5 dB steps in power. This resulted in gain values between 20 and 26 dB, which can also be seen in the variation in the maximum gain of the plotted traces in figure Fig. 4.5 (a). However, the associated gain variations do not appear to be the dominant source of the observed fluctuation in the GBWP. At lower operating frequencies, the GBWP lies slightly below the resonator linewidth $\kappa_{\text{tot}}/2\pi$, whereas at higher frequencies it exceeds this value. The origin of this asymmetry is not yet understood.

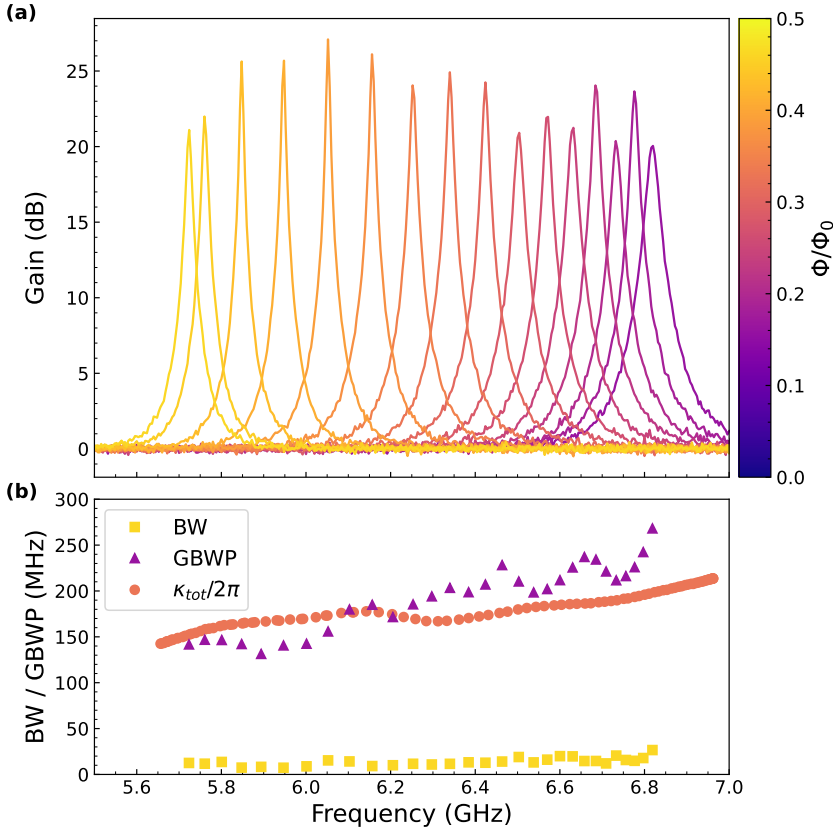


Figure 4.5: Gain, bandwidth, and gain-bandwidth product as a function of the SPA frequency. (a) Measured gain of the SPA across its tunable operating frequency range for a set of flux bias points between $\Phi/\Phi_0 \approx 0.16$ and $\Phi/\Phi_0 \approx 0.46$. Across this range, the device maintains more than 20 dB of gain from roughly 5.7 to 6.8 GHz, corresponding to a tunability of about 1.1 GHz. Each trace corresponds to a different flux bias, indicated by color. (b) Corresponding bandwidth and gain-bandwidth product. The resonator linewidth $\kappa_{\text{tot}}/2\pi$ is included for reference.

GAIN COMPRESSION

Finally, we characterize the dynamic range of the SPA by measuring gain compression as a function of the input signal power. This measurement indicates the maximum input power that the amplifier can handle before its gain is reduced. Figure 4.6 shows the gain compression at a SPA frequency of $\omega_{\text{SPA}}/2\pi = 5.887$ GHz. The pump frequency is set to $\omega_p = 2\omega_{\text{SPA}}$ and the pump power is adjusted to achieve 20 dB of small-signal gain. The 1 dB compression point $P_{-1\text{dB}}$, is defined as the input power at which the gain decreases by 1 dB from its nominal value. For this operating point, we find $P_{-1\text{dB}} \approx -118$ dBm. Where the signal power at the SPA input is calibrated using the method described in supplementary information 4.5.5.

Gain compression can be explained by two main mechanisms [38]. The first mechanics be-

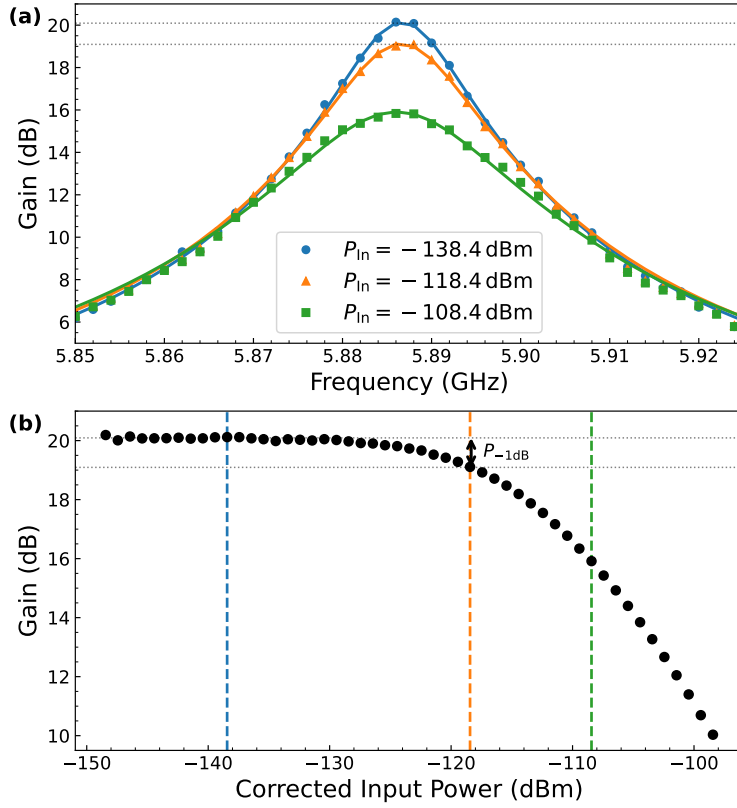


Figure 4.6: Gain compression measurement. (a) Gain as a function of the signal frequency for three different input signal powers. Markers denote experimental data and solid lines are Lorentzian fits. The SPA resonance frequency is $\omega_{\text{SPA}}/2\pi = 5.887$ GHz. (b) The 1 dB compression point $P_{-1\text{dB}}$ is identified as the input power at which the gain is reduced by 1 dB relative to the small-signal gain.

hind gain compression is the Kerr-induced frequency shift. At high input powers, the large population of signal photons in the resonator will induce a Kerr-shift because of the residual self-kerr nonlinearity. When this frequency shift becomes comparable to the SPA linewidth, the pump no longer satisfies the pump condition $\omega_p = 2\omega_{\text{SPA}}$ required for three-wave mixing, and the gain is reduced. The second mechanism behind gain compression is pump depletion. During amplification, pump photons are converted into amplified signal photons. When the input power is small compared to the pump, the effect on the pump amplitude is negligible, and the pump can be considered stiff. However, as the input signal increases and the output power approaches the pump power, a significant fraction of pump photons are converted into amplified signal photons. This breakdown of the stiff-pump approximation leads to pump depletion, reducing the pump amplitude and thus lowering the achievable gain [45].

Analytical estimates for the 1 dB compression power associated with these mechanisms are obtained following the method in Ref. [38]. In the stiff-pump approximation, Kerr-induced gain compression is expected when the nonlinear frequency shift becomes comparable to the SPA

linewidth, yielding

$$P_{-1\text{dB}}^{\text{Stark}} \sim \frac{\kappa}{|g_4|} G_0^{-5/4} \hbar \omega_{\text{SPA}} \kappa. \quad (4.5)$$

Similarly, gain compression due to pump depletion is estimated as

$$P_{-1\text{dB}}^{\text{pump}} \sim \frac{\kappa}{g_3^2 / \omega_{\text{SPA}}} G_0^{-3/2} \hbar \omega_{\text{SPA}} \kappa. \quad (4.6)$$

Using the nonlinear coefficients g_3 and g_4^* extracted from the flux-dependent resonance frequency fit described in Supplementary Information 4.5.6, together with $\kappa/2\pi = 136$ MHz and $G_0 = 20$ dB, we estimate $P_{-1\text{dB}}^{\text{Stark}} \approx -122$ dBm and $P_{-1\text{dB}}^{\text{pump}} \approx -118$ dBm.

These estimates should be regarded as indicative scales rather than sharp bounds, as they rely on indirect extraction of the nonlinear coefficients. The experimentally observed compression point $P_{-1\text{dB}} \approx -118$ dBm lies between these two scales, indicating that both Kerr-induced detuning and pump depletion may contribute to gain compression at this operating point. We note that a more direct experimental extraction of the nonlinear coefficients, for example via pump-power-dependent gain measurements for g_3 and driven Stark-shift measurements for g_4 , would provide a more precise determination of the relative contributions of these two mechanisms.

4.4. CONCLUSION

We presented the design, fabrication, and experimental characterization of a flux-tunable SNAIL-based parametric amplifier (SPA). The devices are based on a $\lambda/2$ microstrip resonator incorporating an array of 20 SNAILS, providing a controllable three-wave-mixing nonlinearity while suppressing unwanted Kerr effects. All devices were fabricated using a reproducible wafer-scale process and characterized at millikelvin temperatures.

The SPA resonator was first characterized in its linear, unpumped regime. The resonance frequency exhibits a periodic flux dependence and is tunable from approximately 5.7 GHz to 7.0 GHz, corresponding to a total tuning range of about 1.3 GHz. Over this range, the total resonator linewidth varies with flux and reaches values on the order of $\kappa_{\text{tot}}/2\pi \sim 200$ MHz near zero flux. Independent measurements on a test resonator confirm that internal dissipation is small compared to the external coupling rate.

When operated as a parametric amplifier, the SPA provides stable gain exceeding 20 dB over an effective operating range from roughly 5.7 GHz to 6.8 GHz, corresponding to a flux-tunable operational bandwidth of approximately 1.1 GHz. At a fixed flux bias point, we observe instantaneous bandwidths between 7 and 27 MHz at 20 dB gain, consistent with the expected gain–bandwidth trade-off. The gain–bandwidth product approaches the total resonator linewidth, in agreement with expectations for resonator-based parametric amplifiers. The dynamic range of the SPA was characterized through gain-compression measurements. At 20 dB of small-signal gain, a 1 dB compression point of $P_{-1\text{dB}} \approx -118$ dBm is obtained at the device input.

The wafer-scale fabrication and performance demonstrated here provide a robust platform for integrating SPAs into complex quantum measurement systems. Continued improvements in dynamic range and bandwidth will directly benefit high-fidelity qubit readout, multiplexed measurement architectures, and broadband microwave optomechanical experiments.

4.5. SUPPLEMENTARY INFORMATION

4.5.1. SPA DESIGN

The SPA devices were designed to fill a 2-inch Silicon wafer hosting 22 individual SPA chips. The wafer design aims to span a broad tunable frequency range between 3 and 10 GHz. Each chip contains a $\lambda/2$ microstrip resonator with a chain of SNAILS providing the flux-tunable nonlinearity. In addition to the amplifier chips, the wafer included two test resonators to characterize the internal quality factor of the superconducting microstrip geometry. Furthermore, arrays of test junctions were located at all four corners of the wafer. Each SPA contains an array of M SNAILS connected in series, giving a total inductance of

$$L_s^{\text{array}}(\Phi) = M L_s(\Phi), \quad (4.7)$$

where $L_s(\Phi)$ is the flux-dependent inductance of a single SNAIL. Every SNAIL consists of one small junction with inductance L_J/α and three large junctions with inductance L_J . At the design stage, we use the relation between the SNAIL inductance at the sweet spot, and the inductance of single junctions estimated from room-temperature resistance measurements using the Ambegaokar–Baratoff relation [46]. Fig 4.7 shows how the effective inductance of a SNAIL can be simplified at the sweet spot. Starting from the full circuit, which consists of two arms in parallel, the inductance is reduced to an equivalent inductance that is given by

$$\frac{1}{L_s(\Phi = 0)} = \frac{1}{3L_J} + \frac{\alpha}{L_J}, \quad (4.8)$$

where α is the ratio between the inductance of the large junction L_J and the small junction L_J/α . The sweet spot inductance of the full SNAIL array can then be related to the large junctions inductance by

$$L_s^{\text{array}}(\Phi = 0) = M L_s(\Phi = 0) = M \frac{L_J}{\alpha + \frac{1}{3}}. \quad (4.9)$$

The SNAIL array is embedded in a $\lambda/2$ microstrip resonator with bare resonance frequency ω_0 and characteristic impedance Z_a . The microstrip geometry uses a non-superconducting

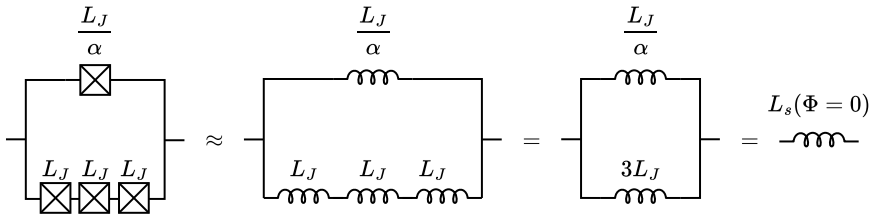


Figure 4.7: Schematic representation of the SNAIL and its equivalent inductive models at the sweet spot. The SNAIL consists of one small junction with inductance L_J/α and three large junctions with inductance L_J . This configuration can be approximated by an equivalent circuit at the sweet spot.

groundplane on the back of the chip. This avoids flux focusing effects allowing for homogeneous flux biasing of all SNAILS by an external superconducting coil. This flux bias controls the nonlinear inductance of the SNAIL array, tuning both the resonance frequency and the nonlinearities. For the fundamental mode $\hat{a}$, the SPA can be described by

$$\mathcal{H}_{\text{SPA}}/\hbar = \omega_{\text{SPA}} \hat{a}^\dagger \hat{a} + g_3 (\hat{a} + \hat{a}^\dagger)^3 + g_4 (\hat{a} + \hat{a}^\dagger)^4, \quad (4.10)$$

where ω_{SPA} is the flux-dependent SPA frequency, g_3 is the cubic coupling (three-wave mixing) set by the SNAIL asymmetry α and the flux bias. g_4 the quartic (Kerr) nonlinearity. The parameters (ω_{SPA} , g_3 , g_4) depend on the device design parameters (L_J , α , M , ω_0 , Z_a) and on the applied external flux Φ .

Fig. 4.8 (a) shows the final wafer layout containing 22 SPA chips. Panel 4.8 (b) highlights the specific chip characterized in this work. This device corresponds to one of the mid-range frequency designs, intended for operation around 6–7 GHz. Its target design parameters are $L_J = 60$ pH, $\alpha = 0.09$, $M = 20$, $\omega_0 = 2\pi \cdot 13$ GHz, and $Z_a = 50$ Ω .

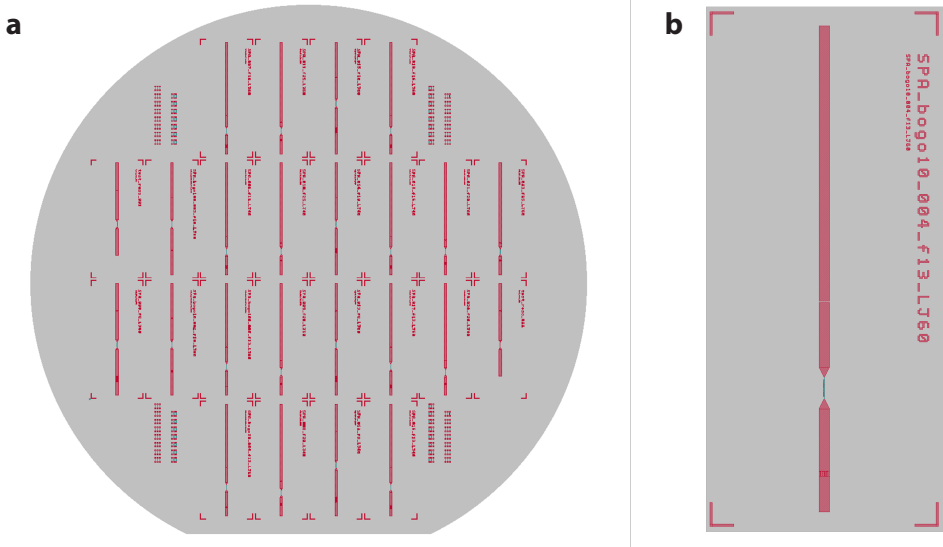


Figure 4.8: SPA wafer layout. (a) Design for fabricating 22 SPA chips over a large frequency range on a 2-inch silicon wafer. (b) Zoom in of a single SPA chip, highlighting the device characterized in this chapter. The $\lambda/2$ microstrip resonator is strongly coupled to the signal port at the bottom and weakly coupled to the pump port at the top.

4.5.2. FABRICATION

Here, we provide an overview of the fabrication process used to realize the SPAs. Additional information on individual fabrication steps and processes can be found in the general fabrication section of Chapter 2.2.

The fabrication starts from a 2-inch double side polished (DSP) high-resistivity silicon wafer. To protect the wafer surface during initial metal deposition, one side is coated with a layer of PMMA A6 spun at 1000 RPM and baked at 180 °C for 5 minutes. The wafer is then loaded into an AJA electron-beam evaporator for deposition of the ground plane required for the microstrip geometry. A 5 nm titanium adhesion layer is first evaporated at a rate of 1 Å/s, followed by 200 nm of gold deposited at 2 Å/s. After metal deposition, the protective resist layer is stripped in hot acetone (50 °C), followed by rinsing in isopropanol (IPA) and a brief (O_2) plasma descum to remove residual organic contamination.

For the Josephson junction fabrication, a multilayer electron-beam resist stack is prepared to enable a clear undercut for shadow evaporation. First, a layer of MMA (8.5) MAA EL13 is spun at 3000 RPM and baked at 180 °C for 3 minutes, followed by a second layer of the same resist spun at 3500 RPM and baked under identical conditions. A top layer of PMMA A4 (950k) is then spun at 1000 RPM and baked at 180 °C for 15 minutes. This results in a resist stack consisting of approximately 1.5 μm of MMA and 300 nm of PMMA. The pattern, shown in Fig. 4.8 (a), is written using a Raith EBPg-5200 electron-beam lithography system. To accommodate the different feature sizes present in the wafer design, the pattern is subdivided into five exposure regions corresponding to the bulk structures and the large and small Josephson junctions. Each junction region is supplemented with a low-dose “ghost” exposure to create a large controlled undercut in the MMA layer during development. The exposed resist is developed in a chilled (7 °C) H_2O :IPA (1:3) solution for 2 minutes, followed by a 15 second rinse in IPA.

After development, the wafer undergoes a short oxygen plasma clean (30 seconds at 200 sccm O_2 , 100 W) to remove any residual resist. The wafer is then transferred to a Plassys electron-beam evaporator for double-angle aluminium evaporation. The first aluminium layer, 30 nm thick, is evaporated at a rate of 0.2 nm/s at an angle of 35°. The wafer is subsequently oxidized in situ at a pressure of 5 mbar for 5 minutes to form the tunnel barrier. After a 180° rotation of the sample, a second aluminium layer of 100 nm thickness is evaporated at the same rate and angle to form the Josephson junctions.

Liftoff is performed by immersing the wafer in a hot acetone bath (50 °C) for several hours. Remaining metal flakes are gently removed using a pipette, followed by medium-power sonication for 2 minutes. The wafer is then rinsed in IPA and dried with nitrogen. Following liftoff, test junctions are characterized at room temperature to verify junction resistance and confirm that the devices meet the design specifications. To protect the wafer during dicing, a layer of PMMA A6 is spun at 1000 RPM and baked at 100 °C for 1 minute. The wafer is then diced into individual chips of size 11 mm $\times$ 5 mm. After dicing, the protective resist is stripped in hot acetone (50 °C) for 30 seconds, followed by an IPA rinse and nitrogen drying. A final oxygen plasma descum (3 minutes at 200 sccm O_2 , 100 W) is performed to remove residual contamination, after which the SPA chips are ready for packaging and measurement.

4.5.3. MEASUREMENT SETUP

The experiments were carried out with the SPA mounted on the base plate of a Bluefors dilution refrigerator operating at base temperature ($T \approx 25$ mK). A schematic of the wiring within the fridge is shown in Fig. 4.9. The signal input line has 40 dB of room-temperature attenuation and 53 dB of cryogenic attenuation. The input signal then passes through a 4-8 GHz circulator (LNF-CIC4_8A) before reaching the signal port of the SPA where it is reflected. The reflected signal returns to the circulator and passes through an infrared filter, Low-pass filter (ZLSS-11G-S+) and two isolators before being further amplified by a high electron mobility transistor (HEMT) amplifier (Low Noise Factory LNF-LNC4_8C) and a room temperature amplifier (Narda-MITEQ). The pump port of the SPA is connected to a separate pump line that has 33 dB of cryogenic

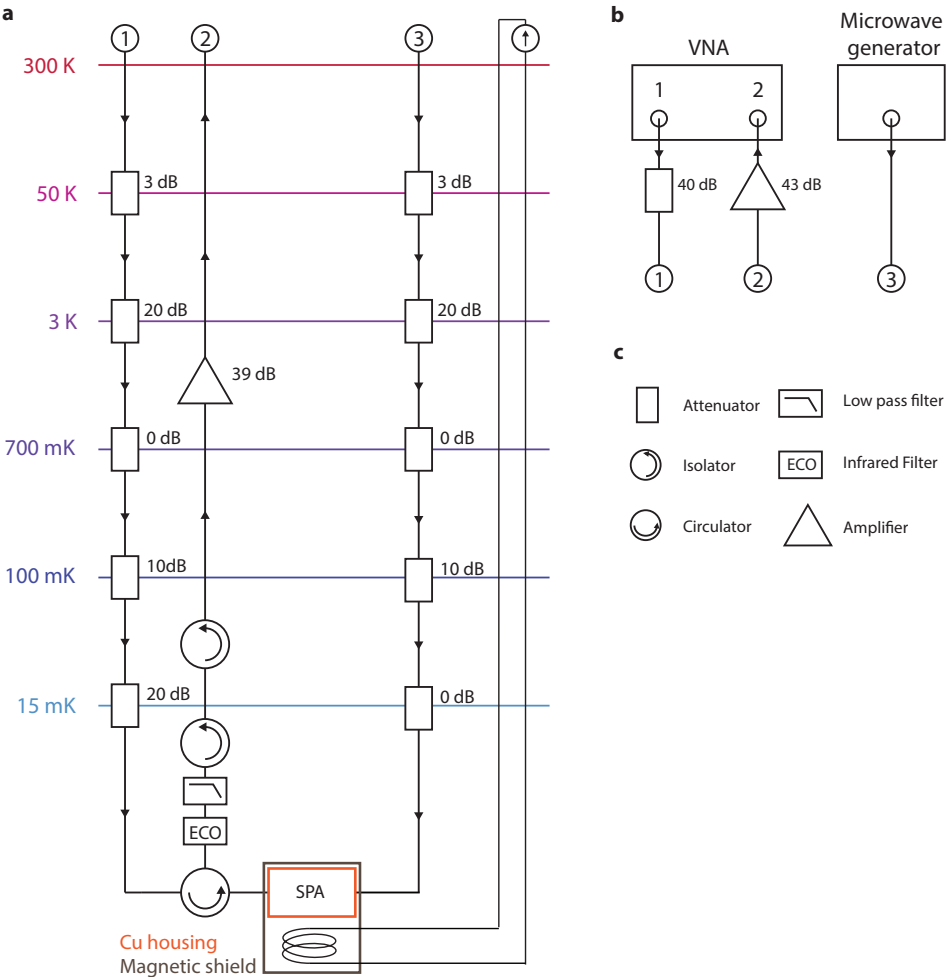


Figure 4.9: Schematic of the measurement setup. (a) Wiring diagram inside the dilution refrigerator. **(b)** Room temperature setup for VNA measurements. **(c)** Overview of components.

attenuation and no extra room temperature attenuation. The pump tone is generated by a Keysight E8257D signal generator. A superconducting NbTi coil positioned at the back of the PCB within the Mu-metal magnetic shielding box of the SPA packaging applies the magnetic flux to the SNAILs. The coil is connected to the DC-loom which connects to the current source (Stanford RS CS580) at room temperature. Microwave measurement are performed using a keysight N522A PNA network analyzer.

4.5.4. MAGNITUDE BACKGROUND CORRECTION

The raw magnitude data of the linear reflection measurement is shown in Fig. 4.10 (a). Slow variations in the background can be observed in the line trace prior to the background correction that is shown in Fig. 4.2 (b), which are expected to originate from frequency dependent attenuation in the microwave components and cabling. To isolate the response of the SPA, we applied the following background subtraction procedure.

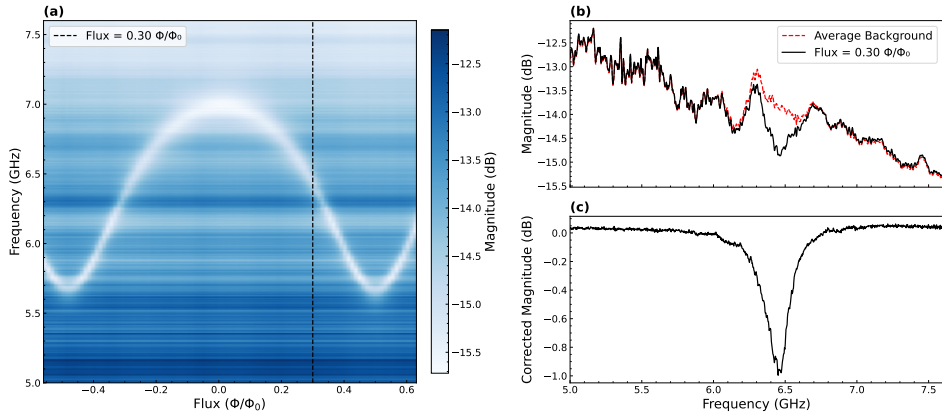


Figure 4.10: Background subtraction of the reflection magnitude. (a) Raw magnitude response of the SPA as a function of probe frequency and applied magnetic flux. This dataset corresponds to the measurement shown in Fig. 4.2(b) before background correction. The dashed line indicates the flux bias point at $\Phi/\Phi_0 = 0.30$, where the line cut shown in panel (b) is taken. **(b)** Magnitude trace at $\Phi/\Phi_0 = 0.30$. The red dashed line represents the averaged background obtained by averaging the off-resonant regions of all flux traces. **(c)** Corrected magnitude response after subtraction of the averaged background. This processed data is used to produce the main-text result in Fig. 4.2(d).

For each flux bias point, the complex trace S_{21} was fitted to determine the resonance frequency $\omega_{\text{SPA}}/2\pi$. A frequency window of ± 0.5 GHz centered around $\omega_{\text{SPA}}/2\pi$ was excluded to remove the resonance from the background estimation. This frequency window was chosen to be more than 5 times the resonance linewidth. The remaining parts of the off-resonant traces were added together and averaged to obtain a smooth, flux independent background function $B(\omega)$. Finally, this averaged background was subtracted from each individual trace according to

$$|S_{21}^{\text{corr}}(\omega, \Phi)| = |S_{21}(\omega, \Phi)| - B(\omega), \quad (4.11)$$

where $|S_{21}^{corr}(\omega, \Phi)|$ denotes the corrected magnitude shown in Fig. 4.2 (a). This procedure effectively removed broadband transmission variations while preserving the resonance amplitude response of the SPA. Fig. 4.2 (b) shows a magnitude response trace at a single flux point $0.30 \Phi / \Phi_0$. In the original magnitude trace, there is no clear resonant feature. But the difference becomes apparent when we also plot the average background determined using all other flux points.

As can be seen in Fig. 4.2 (c), the resulting corrected data exhibits a flat baseline and now clearly shows the residual magnitude feature associated with the resonant reflection. We note that this correction is still not perfect, as can be seen in the feature that is left around 6.1 GHz. The reason for this is that the applied correction only accounts for frequency dependent attenuation of the signal path. It does not account for interference between different signal paths [41]. Eliminating these effects would require a more elaborate short-open-load (SOL) calibration, which characterizes both the amplitude and phase errors.

4.5.5. INPUT POWER CALIBRATION

To estimate the input power at both the pump and signal ports, we calibrate the attenuation of the input lines and include the specified attenuation of additional attenuators in the setup. The line attenuation was measured at a base temperature of $T = 25$ mK by connecting two input lines at the mixing chamber stage with a short coaxial cable and measuring the round-trip transmission. Figure 4.11 shows the resulting transmission data obtained for all input lines in the dilution refrigerator. From a polynomial fit to the measured round-trip transmission, we determine the frequency-dependent attenuation of a single line by dividing the round-trip value by two. This approach accounts for the attenuation of all components in the signal path, including the room-temperature cables that connect to the measurement instruments. The estimated attenuation may be slightly overestimated due to the additional loss introduced by the short cable used to connect the two input lines at the MXC stage. The estimated power at the device is calculated by subtracting the total attenuation from the source power

$$P_{dev}(\omega) \approx P_{src}(\omega) - A_{line}(\omega) - A_{att}, \quad (4.12)$$

where $P_{src}(\omega)$ is the generator output power, A_{line} is the measured frequency-dependent line attenuation, and A_{att} is the fixed attenuation from additional attenuators.

4.5.6. FIT PROCEDURE

To extract the SPA parameters from the measured resonance frequency ω as a function of flux, we use a two-step fitting approach based on the theoretical model described in Ref. [39]. The resonance frequency of the SPA can be described as

$$\frac{2Z_a}{ML_s(\Phi)} = \omega_{SPA} \tan\left(\frac{\pi}{2} \frac{\omega_{SPA}}{\omega_0}\right), \quad (4.13)$$

where Z_a is the characteristic impedance of the microstrip geometry, M is the number of SNAILS, $L_s(\Phi)$ is the flux-dependent inductance of a single SNAIL, and ω_0 is the fundamental

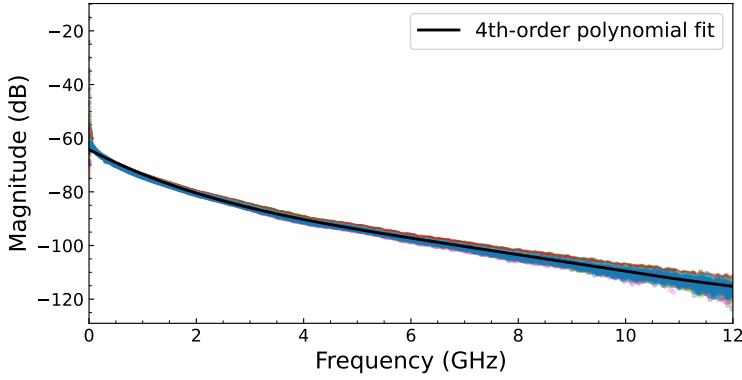


Figure 4.11: Estimation of the frequency-dependent input line attenuation Transmission of two connected input lines measured at $T = 25$ mK. The plotted data shows the frequency-dependent attenuation for all input lines, along with a polynomial fit to obtain the average attenuation at a certain frequency. The single line attenuation is obtained by dividing the obtained round-trip value from the fit by two.

4

frequency of a square microstrip resonator of the same size without the SNAIL array. This equation is solved numerically by evaluating the cost function

$$f(\omega) = \omega_{\text{SPA}} \tan\left(\frac{\pi}{2} \frac{\omega_{\text{SPA}}}{\omega_0}\right) - \frac{2Z_a}{ML_s(\Phi)}, \quad (4.14)$$

over a fine frequency grid that determines the points where the cost function is minimized. We fix the number of SNAILS to $M = 20$ and set the impedance to $Z_a = 50 \Omega$ for our microstrip geometry. The fitting route determines the design parameters (L_J, α, ω_0) and two linear coefficients a and b that are used to map the applied current to the external flux via $\Phi = aI + b$. We first perform a global search using a differential evolution method, followed by local refinements using a least squares method. This approach provides a frequency fit across the entire flux range.

After determining the design parameters from the frequency fit, we evaluate the nonlinear response of the SNAIL array following Ref. [39]. The dimensionless SNAIL potential,

$$U(\varphi) = -\alpha \cos \varphi - 3 \cos\left(\frac{\varphi_{\text{ext}} - \varphi}{3}\right), \quad (4.15)$$

is expanded around its flux-dependent minimum φ_0 in terms of small phase fluctuations $\delta\varphi = \varphi - \varphi_0$, such that

$$U(\varphi_0 + \delta\varphi) = \sum_{n \geq 2} \frac{c_n(\Phi)}{n!} (\delta\varphi)^n, \quad (4.16)$$

with

$$c_n(\Phi) = \left. \frac{\partial^n U(\varphi)}{\partial \varphi^n} \right|_{\varphi=\varphi_0}, \quad c_1 = 0. \quad (4.17)$$

The quadratic coefficient c_2 determines the linearized SNAIL inductance via $L_s = L_J/c_2$, while the higher-order coefficients c_3 and c_4 set the cubic and quartic nonlinearities, respectively. Using these coefficients and the fitted resonance frequency $\omega_{\text{spa}}(\Phi)$, the effective three- and four-wave mixing strengths g_3 and g_4^* are given by

$$g_3 = \frac{4Z_a c_3}{3M^2 L_J} \sqrt{\frac{Z_a}{Z_Q}} \left[\frac{\cos^2\left(\frac{\pi}{2} \frac{\omega_{\text{spa}}}{\omega_0}\right)}{\pi \frac{\omega_{\text{spa}}}{\omega_0} + \sin\left(\pi \frac{\omega_{\text{spa}}}{\omega_0}\right)} \right]^{3/2}, \quad (4.18)$$

$$g_4^* = \frac{1}{12} \frac{\omega_{\text{spa}} \sin^2\left(\pi \frac{\omega_{\text{spa}}}{\omega_0}\right)}{c_2 M^2 \tan\left(\frac{\pi}{2} \frac{\omega_{\text{spa}}}{\omega_0}\right) \left[\pi \frac{\omega_{\text{spa}}}{\omega_0} + \sin\left(\pi \frac{\omega_{\text{spa}}}{\omega_0}\right)\right]^2} \frac{Z_a}{Z_Q} \left[c_4 - \frac{c_3^2}{c_2} \frac{3 + 5\left(\frac{\omega_{\text{spa}} M L_s}{2 Z_a}\right)^2}{1 + 3\left(\frac{\omega_{\text{spa}} M L_s}{2 Z_a}\right)^2} \right]. \quad (4.19)$$

Here g_4^* denotes the effective quartic nonlinearity after a renormalization by the cubic interaction given by

$$g_4^* = g_4 - \frac{5g_3^2}{\omega_a}. \quad (4.20)$$

The observable self-Kerr coefficient is given by $K = 12g_4^*$.

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5

IMPACT OF SURFACE CLEANING ON LOSS IN TANTALUM COPLANAR WAVEGUIDE RESONATORS

Energy dissipation in superconducting microwave circuits limits coherence times and remains a key challenge for superconducting quantum devices. In this work, we study loss mechanisms in α -phase tantalum coplanar-waveguide (CPW) resonators, focusing on surface-related dielectric loss in the single-photon regime. Using resonators fabricated from a single tantalum film with systematically varied surface participation ratios, we investigate the combined effects of resonator geometry and post-fabrication surface treatments. Microwave transmission measurements at millikelvin temperatures show that additional surface cleaning, particularly a combined piranha and buffered oxide etch, significantly reduces loss and enables reproducible single-photon internal quality factors exceeding one million. These results highlight the importance of surface engineering for achieving low-loss tantalum superconducting circuits.

5.1. INTRODUCTION

Energy loss in superconducting microwave circuits remains one of the primary obstacles to achieving long coherence times in quantum devices. Dissipation limits the performance of superconducting resonators and qubits by constraining coherence times, measurement fidelity, and overall device quality. Understanding and mitigating loss mechanisms in superconducting circuits is therefore a key challenge in the development of scalable and high-coherence quantum hardware [1–3].

Superconducting coplanar-waveguide (CPW) resonators are widely used as test structures for investigating loss mechanisms in superconducting circuits [4–9]. Compared to qubits, CPW resonators are simple to fabricate and characterize, enabling rapid design iterations and the systematic exploration of material and geometric parameters. In addition, resonators can be measured from the single-photon regime up to high photon numbers, allowing power-dependent loss mechanisms to be disentangled [10, 11]. While resonators do not capture additional loss channels associated with Josephson junctions, they provide a controlled and versatile platform for isolating linear dissipation mechanisms that are also relevant to superconducting qubits.

At millikelvin temperatures and in the single-photon regime, loss in superconducting resonators is commonly attributed to two-level systems (TLS) residing in amorphous dielectric materials [1, 2, 4]. These TLS are believed to originate primarily from thin dielectric layers at material interfaces, including the metal–substrate (MS), substrate–air (SA), and metal–air (MA) interfaces [12–14]. The contribution of these interfaces to the total loss depends on the fraction of the electric field energy stored near each interface. As a result, the geometry of a CPW resonator enables control of loss through the surface participation ratio (SPR), which quantifies the fraction of electric field energy stored in lossy surface regions [15–17].

Recently, tantalum has emerged as a promising superconducting material for high-coherence quantum devices [18–20]. Superconducting qubits fabricated from tantalum have demonstrated exceptionally long coherence times, motivating increased interest in this material system. These improvements have been attributed to the chemical stability of tantalum and the formation of a thin, robust native oxide, which may host fewer lossy TLS compared to other commonly used materials such as aluminum [21–23]. Recent work has demonstrated that the careful engineering of resonator geometry and surface treatments can significantly reduce dielectric loss and enable internal quality factors exceeding one million in planar superconducting resonators [24, 25].

In this chapter, we present a systematic study of loss in tantalum CPW resonators. Using devices fabricated from a single tantalum film, we investigate how resonator geometry and post-fabrication cleaning processes affect the internal quality factor in the single-photon regime. By comparing resonators with different surface participation ratios and surface treatments, we identify fabrication strategies that enable internal quality factors to exceed one million in the single-photon regime.

5.2. DEVICE DESIGN AND FABRICATION

To systematically investigate how post-fabrication cleaning processes can reduce dielectric losses in superconducting resonators, we design a chip containing twelve quarter-wavelength ($\lambda/4$) coplanar waveguide (CPW) resonators capacitively coupled to a common feedline. This provides a controlled platform in which resonator loss can be compared across different device geometries and surface treatments. An overview of the full chip layout is shown in Fig. 5.1 (a). Each resonator is shorted to ground at one end and open at the other, where it couples capacitively to the central feedline, as illustrated in the zoom-in shown in Fig. 5.1 (b). The resonators are designed to span a frequency range between 5.0 and 7.2 GHz, with an average frequency spacing of approximately 200 MHz. This spacing is chosen to ensure that individual resonances are well separated and can be reliably identified in transmission measurements.

To disentangle the influence of geometry, coupling strength, and surface loss mechanisms, the twelve resonators are fabricated with two distinct target coupling rates and three different surface participation ratios (SPRs). The coupling rates are alternated across the frequency-ordered resonators, such that each measured resonance can be uniquely mapped back to its corresponding design. The presence of undercoupled resonators is important for a reliable extraction of the internal quality factor Q_i . Overcoupled resonators are prone to large uncertainties in Q_i due to Fano interference [26], which leads to systematic errors in the extraction of internal losses in these hanger-type transmission measurements. The three SPR designs are realized by varying the CPW center conductor width w and gap width g , as summarized

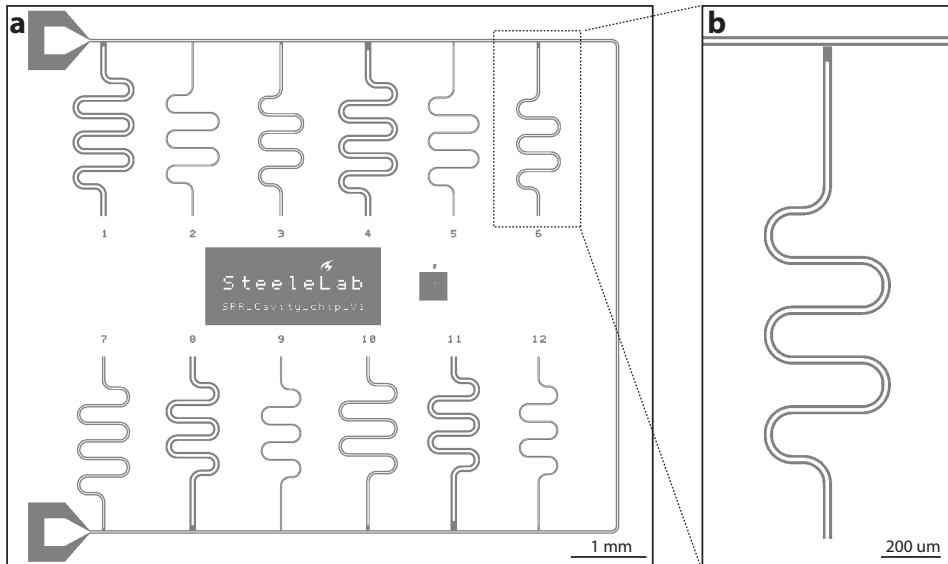


Figure 5.1: Overview of the full chip layout. (a) Twelve quarter-wavelength ($\lambda/4$) coplanar-waveguide resonators are coupled to a main feedline. The resonators are designed with three different surface participation ratios and two different target coupling rates. The large exposed square in the center of the chip is used for end-point detection of the etching process. **(b)** The zoom-in of a single resonator shows how the resonators are capacitively coupled to the main feedline.

in SI Table 5.1. By modifying these dimensions, the fraction of the electric field energy stored in lossy surface layers, such as the metal–substrate (MS), substrate–air (SA), and metal–air (MA) interfaces is systematically controlled. This enables a comparison of resonator loss as a function of surface participation while keeping all other device parameters fixed.

All devices presented in this chapter were fabricated on 4-inch double-side polished (DSP) C-plane sapphire (Al_2O_3) wafers with a thickness of 430 μm . The superconducting film consists of 100 nm of α -phase tantalum deposited by STAR Cryoelectronics. The wafer was diced into 10 mm $\times$ 10 mm chips, which we use to fabricate the devices measured in this work. To eliminate variations originating from differences in film deposition, all resonators were fabricated from this same wafer. Patterning of the resonators was performed using electron-beam lithography followed by metal etching in a Sentech Etchlab 200 reactive-ion etching (RIE) system using CF_4 . Endpoint detection was performed on a large exposed region in the center of the chip to determine when the tantalum film had been fully etched away, after which the etch was continued for an additional 10 seconds to ensure that the film was completely removed across the entire chip. CF_4 based etching selectively removes tantalum while leaving the sapphire substrate mostly unaffected. After etching, the resist was stripped and the chip was cleaned using an O_2 plasma descum. This process defines the reference fabrication flow used for all devices in this chapter. Further fabrication details are provided in Supplementary Information 5.5.2. In total, seven chips were measured, corresponding to 83 individual resonators with varying geometries and surface treatments.

5

5.2.1. POST-FABRICATION SURFACE TREATMENTS

Following the reference fabrication process described above, additional post-fabrication cleaning steps were applied to evaluate how different fabrication treatments impact resonator performance, with the goal of minimizing loss in tantalum CPW resonators. These treatments were chosen to target different classes of surface contamination, including native oxides and organic residues that are expected to contribute to dielectric loss.

These treatments were selected based on fabrication and surface-cleaning approaches used in prior studies of tantalum superconducting resonators and qubits, where similar processes have been shown to reduce dielectric loss [18, 20].

Four distinct processing conditions were studied:

- **Reference process:** No additional cleaning beyond resist stripping and an O_2 plasma descum, serving as a baseline fabrication process.
- **Buffered oxide etch (BOE):** Immersion in a 7:1 buffered HF solution for 20 minutes to remove native surface oxides.
- **Piranha clean:** A 20 minute piranha clean ($\text{H}_2\text{SO}_4\text{:H}_2\text{O}_2$) to remove organic contaminants and resist residues.
- **Combined piranha and BOE:** Sequential piranha and BOE cleaning, each applied for 20 minutes, to remove both organic residues and native oxides.

After the cleaning steps, the chips were rinsed in water, dried with nitrogen, and packaged. Because all devices originate from the same wafer and undergo identical fabrication steps prior to the cleaning step, observed variations in resonator performance are therefore primarily attributed to differences in surface treatment. We fabricated and measured 7 chips with a total of 83 observed resonators with varying geometries and cleaning procedures. For each cleaning condition, the cleaning procedure was applied simultaneously to two chips at the same time to assess the chip-to-chip variability and verify the reproducibility of the observed resonator performance. The BOE cleaning condition was only evaluated using a single chip.

5.3. MICROWAVE CHARACTERIZATION AND LOSS EXTRACTION

The loss properties of the resonators were characterized by measuring the complex transmission coefficient S_{21} of the common feedline using a vector network analyzer (VNA). All measurements were performed in a Bluefors dilution refrigerator operating at a base temperature of approximately 15 mK. The wiring diagram of the setup can be found in Supplementary Information Fig. 5.10. Each chip contains twelve unique $\lambda/4$ CPW resonators coupled to the same feedline, allowing all resonators on the chip to be probed simultaneously under identical experimental conditions.

Figure 5.2 shows an example of the measured transmission response for a single resonator. Panels (a) and (b) display the magnitude and phase of S_{21} as a function of frequency, while panel (c) shows the corresponding trajectory in the complex plane. In the raw transmission data, the resonant response is superimposed on a slowly varying background in both amplitude and phase. This background originates from frequency-dependent attenuation and phase shifts in the microwave components and cabling of the measurement chain. To remove this effect, we performed a complex background correction. The correction removes linear frequency-dependent variations in the real and imaginary components of S_{21} , and was applied separately for each different measurement configuration. A detailed description of this procedure is provided in Supplementary Information 5.5.4.

Each resonance was then fitted at each probe power using a hanger-type resonator model. A large variety of models have been developed to describe microwave resonator transmission measurements [11, 26–29]. Here, we adopt the model introduced in Ref. [29]. From these fits, we extract the resonance frequency $\omega_0/2\pi$, the internal quality factor Q_i , and the coupling quality factor Q_c . The complex transmission coefficient is modeled as

$$S_{21}(\omega) = a_{in} e^{i(\omega - \omega_0)T} \left[1 - \frac{\frac{Q_{tot}}{Q_c} + 2i \frac{Q_{tot}\alpha}{\omega_0}}{1 + 2i Q_{tot} \left(\frac{\omega - \omega_0}{\omega_0} \right)} \right] \quad (5.1)$$

where ω_0 is the resonance frequency, Q_c and Q_{tot} are the coupling and total quality factors, α is an asymmetry parameter that accounts for Fano-type interference in the measurement chain, for example arising from parallel signal paths or reflections caused by impedance mismatches. The complex prefactor A_{in} captures an overall amplitude and phase offset, while T accounts for the electrical delay. The internal quality factor is obtained from $\frac{1}{Q_i} = \frac{1}{Q_{tot}} - \frac{1}{Q_c}$.

The applied microwave drive power is converted to an average resonator photon number $\langle n \rangle$ using the fitted resonance parameters and the calibrated input-line attenuation. Details of this photon-number calibration procedure are given in Appendix 5.5.6. This enables a direct comparison of resonator properties as a function of photon number across different devices and cleaning processes. The data shown in Fig. 5.2 correspond to a drive power yielding an average photon number $\langle n \rangle \approx 1$, placing it in the single-photon regime. At this power, the extracted internal quality factor is $Q_i = (1.18 \pm 0.04) \times 10^6$. This demonstrates that we can obtain single-photon internal quality factors exceeding one million for resonators with low surface participation ratios fabricated using the combined Piranha and BOE cleaning process.

In the remainder of our analysis, we restrict ourselves to undercoupled resonators ($Q_i < Q_c$). Only including undercoupled resonators provides a more robust metric for comparing the resonator Q_i between the different geometries and surface cleaning procedures.

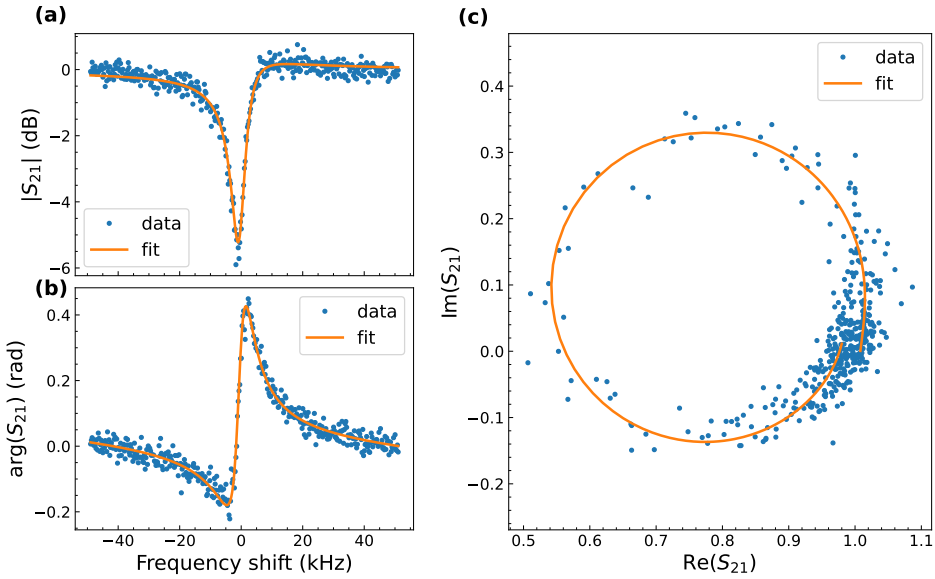


Figure 5.2: Single-photon resonator response with an internal quality factor exceeding one million. (a) Magnitude of the transmission response $|S_{21}|$ of a linear microwave resonator with low surface participation ratio fabricated using a Piranha + BOE cleaning process. **(b)** Phase of the complex transmission S_{21} near resonance. **(c)** Complex S_{21} data in the IQ plane. In (a) and (b), the horizontal axis is the frequency shift Δf relative to the resonance frequency $\omega_0/2\pi = 5.169$ GHz. Circles denote measured data and solid lines show fits to the data. All data shown have their background corrected. The trace shown corresponds to a drive power yielding an average photon number $\langle n \rangle \approx 1$, for which the internal quality factor is $Q_i = (1.18 \pm 0.04) \times 10^6$, demonstrating single-photon operation with internal quality factors exceeding one million.

PHOTON-NUMBER DEPENDENCE OF RESONATOR LOSS

To further investigate the loss mechanisms in the resonators, we study the dependence of the resonance frequency and quality factors as a function of the average photon number. Fig. 5.3 (a) and (b) show the power dependence of the resonance frequency and quality factors for a resonator with high surface participation ratio fabricated using the reference process without additional cleaning. At low photon numbers, the resonance frequency remains approximately constant. However, a downward shift of the resonance frequency is observed at high powers. This frequency shift is accompanied by a reduction of Q_i and is consistent with microwave-induced heating and the generation of nonequilibrium quasiparticles at these elevated drive powers [30]. The extracted coupling quality factor Q_c remains constant over the entire power range. This is expected as it is set by the fixed capacitive coupling to the feedline. In contrast, the internal quality factor Q_i exhibits a clear dependence on photon number. At the lowest driver powers ($\langle n \rangle \leq 1$), it saturates to a constant value, indicating a regime in which the dominant two-level systems (TLS) remain unsaturated. As the photon number is increased, Q_i rises due to the progressive saturation of TLS in the dielectric interfaces. At sufficiently high photon numbers, the internal quality factor eventually degrades again, consistent with the onset of additional non-TLS loss mechanisms, such as quasiparticle-related dissipation associated with heating.

Figure 5.3 (c) compares the internal quality factor as a function of photon number for lithographically identical resonators with different post-fabrication cleaning processes. All devices show a broadly similar power dependence, with an increase in Q_i as the photon number is increased, consistent with the progressive saturation of two-level systems (TLS). However, the absolute value of Q_i at a given photon number depends on the applied cleaning process, indicating that surface treatment has a direct impact on resonator loss. Notably, for some resonators, the lowest accessible photon numbers do not reveal a well-defined saturation plateau. As a result, a full quantitative fit to a TLS loss model is not possible for all devices, since the unsaturated TLS regime is not always observed within the available power range. To enable a consistent comparison across different resonator geometries and surface treatments, we therefore use the internal quality factor measured near the single-photon level as a comparison metric. Although this value does not necessarily correspond to a fully TLS-limited regime, this quantity can be reliably extracted for all measured devices and is directly relevant for typical operating conditions in circuit QED experiments.

The black crosses in Fig. 5.3 (c) indicate the internal quality factor evaluated at ($\langle n \rangle \leq 1$). From this, we observe that resonators subjected to additional post-fabrication cleaning steps exhibit systematically higher single-photon internal quality factors compared to devices fabricated using the reference process. To quantify this trend and disentangle the effects of surface treatment from those of resonator geometry, we next compare the single-photon internal quality factors across all measured devices as a function of surface participation ratio and cleaning procedure.

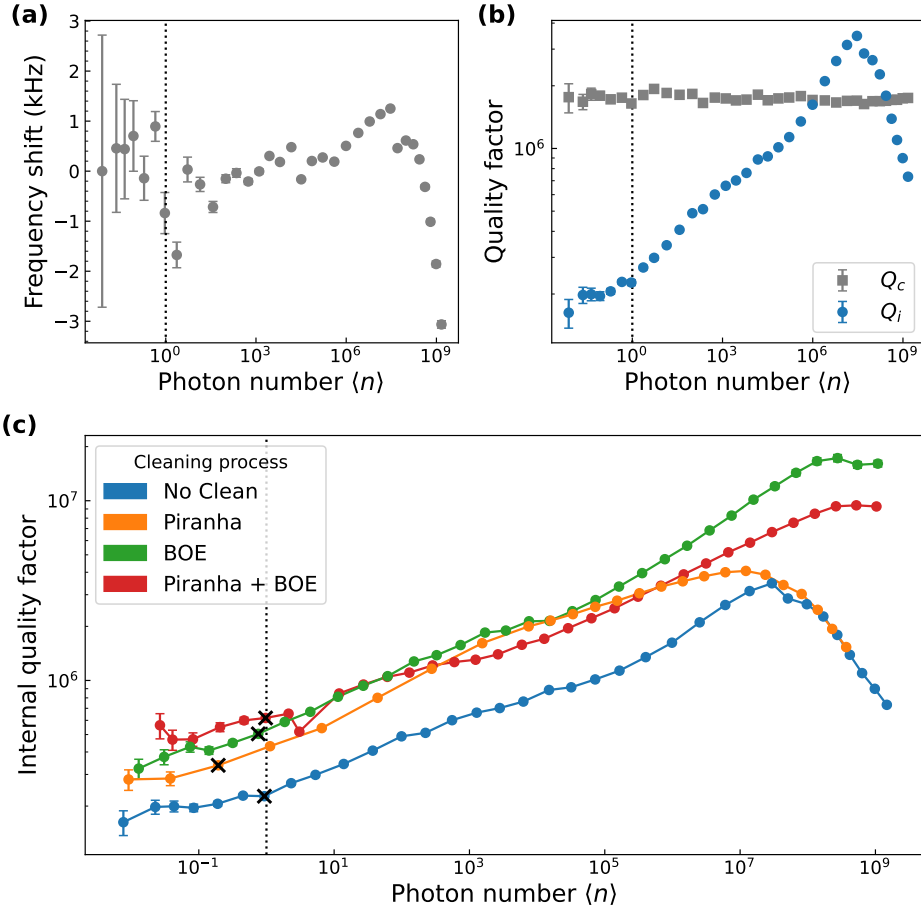


Figure 5.3: Extracted resonator properties as a function of photon number. (a) Frequency shift Δf as a function of the average photon number $\langle n \rangle$ for a high surface participation ratio resonator without additional cleaning. The shift is shown relative to the resonance frequency $\omega_0/2\pi = 5.605$ GHz. (b) Internal (Q_i) and external (Q_c) quality factors extracted from the same resonator as a function of photon number. The external quality factor is approximately constant while the internal quality factor increases with photon number. (c) Internal quality factor Q_i versus photon number for lithographically identical resonators fabricated using different cleaning processes. Black crosses indicate the single-photon internal quality factors. These values are used for the comparison of cleaning methods in subsequent figures.

SURFACE PARTICIPATION DEPENDENCE AND THE EFFECT OF CLEANING

Fig. 5.4 summarizes the single-photon internal quality factors extracted from all the under-coupled resonators measured across seven different chips, grouped by surface participation ratio (SPR) and post-fabrication cleaning procedure. A clear dependence on the surface participation ratio is observed for all different cleaning conditions. Resonators with higher SPR systematically show lower internal quality factors, consistent with an increased sensitivity to surface-related dielectric losses. This behavior supports the interpretation that surface losses dominate the dissipation in these devices and validates the use of SPR as a control parameter to investigate surface loss effects.

Within each SPR group, the impact of post-fabrication cleaning can be clearly observed. High-SPR resonators, which are the most sensitive to surface loss, show the clearest improvement with additional cleaning steps. Applying an additional cleaning step, either a 20-minute Piranha clean or 20-minute buffered oxide etching (BOE), improves the single-photon internal quality factor relative to the reference process. Among the individual treatments, BOE yields a

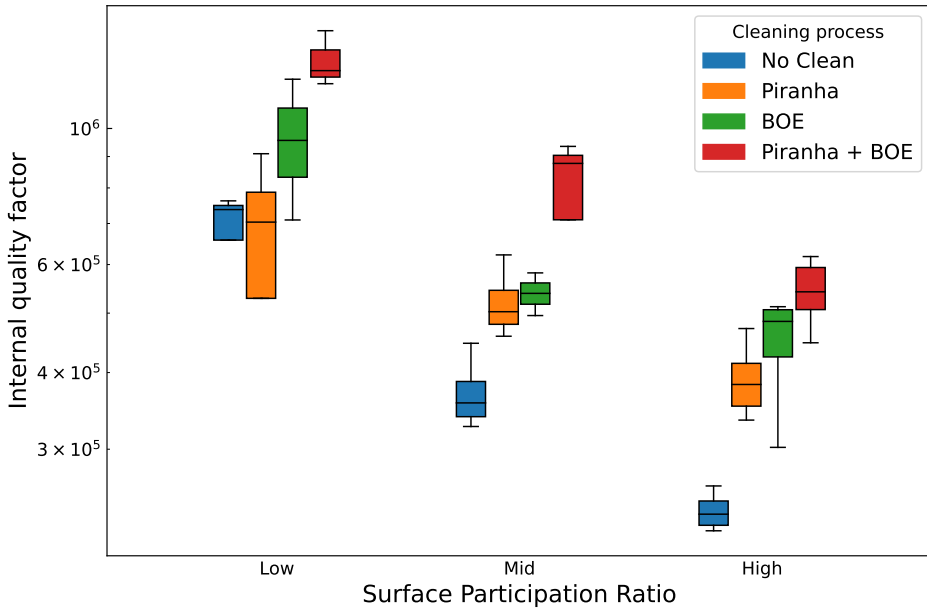


Figure 5.4: Internal quality factor as a function of surface participation ratio and cleaning process. Single-photon internal quality factors Q_i extracted at $\langle n \rangle \leq 1$ are shown for resonators with low, intermediate, and high surface participation ratios. Only undercoupled resonators satisfying $Q_i < Q_c$ are included. The results are summarized using box plots that capture the distribution of single-photon Q_i values across all chips for each geometry and cleaning procedure. Boxes indicate the interquartile range (25th–75th percentile), with the median shown as a horizontal line. Whiskers extend to 1.5 times the interquartile range and outliers are not shown. Resonators with high surface participation ratio are most sensitive to surface-related losses and therefore provide the clearest comparison between different cleaning methods. For low surface participation ratio devices, the combined Piranha + BOE cleaning process reproducibly yields internal quality factors exceeding 10^6 .

larger improvements compared to Piranha. However, the highest single-photon internal quality factors are achieved when a Piranha clean is followed by BOE, demonstrating the benefit of combining both processes.

For low-SPR resonators, the dependence on cleaning is less pronounced, consistent with a reduced contribution of surface-related losses. Nevertheless, a systematic upward shift of the Q_i distribution is still observed. In particular, the combined Piranha + BOE process reproducibly yields single-photon internal quality factors that exceed 10^6 , demonstrating that post-fabrication surface cleaning can play an important role in reducing loss in tantalum superconducting resonators.

5.4. CONCLUSION AND OUTLOOK

In this chapter, we investigated loss mechanisms in tantalum coplanar-waveguide resonators, with a particular focus on the role of post-fabrication surface treatments. By fabricating all devices from a single tantalum film and systematically varying both resonator geometry and the surface cleaning procedure, we were able to directly compare resonator performance under well-controlled conditions.

We find that post-fabrication surface cleaning has a clear and reproducible impact on the internal quality factor in the single-photon regime. Resonators subjected to additional cleaning steps consistently exhibit higher Q_i values compared to devices fabricated using the reference process. Among the treatments studied, the combination of a Piranha clean followed by buffered oxide etching yields the largest improvement, enabling single-photon internal quality factors exceeding 10^6 for low surface participation ratio resonators. These results indicate that surface contamination and native oxides at dielectric interfaces indeed play a significant role in limiting the coherence of the resonator in tantalum-based devices.

5.5. SUPPLEMENTARY INFORMATION

5.5.1. DEVICE DESIGN

The chip layout contains two microwave launch pads connected by a coplanar waveguide (CPW) feedline. The feedline has a center-conductor width $w = 20\ \mu\text{m}$ and gap width $g = 12\ \mu\text{m}$, corresponding to a characteristic impedance of approximately $50\ \Omega$. Twelve quarter-wavelength CPW resonators are capacitively coupled to the feedline along its length. The surface participation ratio is tuned through the resonator geometry by varying the CPW center-conductor and gap dimensions. Each chip includes three SPR classes (low, mid, and high), implemented through different resonator geometries. Larger CPW dimensions reduce the electric-field concentration near lossy surface and interface layers, while smaller dimensions increase it. Figure 5.10(a) shows a schematic cross section of the CPW geometry, highlighting the metal–substrate (MS), substrate–air (SA), and metal–air (MA) interfaces. Increasing the CPW dimensions produces a more distributed electric-field profile, with a larger fraction of the energy stored in the bulk substrate and vacuum.

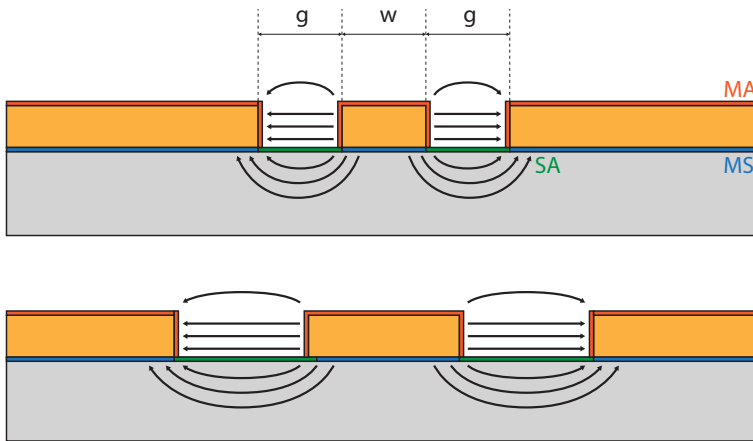


Figure 5.5: Surface participation ratio in a CPW resonator. Schematic cross section of a CPW resonator illustrating how the surface participation ratio depends on device dimensions. The gap g is varied to tune the electric-field distribution, while the trace width w is adjusted to keep the characteristic impedance approximately $50\ \Omega$. As g increases, a smaller fraction of the electric-field energy resides in dissipative surface and interface layers.

Table 5.1 summarizes the full set of design parameters for the twelve resonators on a single chip. Each resonator is assigned an on-chip label corresponding to its identifier in the GDS layout. This labeling differs from the ordering used in the measurement data, which is sorted by increasing resonance frequency. The table therefore enables direct comparisons between resonators with different surface participation ratios and coupling strengths. In support of open science, the resonator chip GDS layout is publicly available on Zenodo [31], facilitating reproducibility and reuse.

Table 5.1: Design parameters for the twelve quarter-wavelength CPW resonators on a single chip. The on-chip label refers to the resonator index in the GDS layout and differs from the ordering used in the measurement data, which is sorted by increasing resonance frequency. The resonators span three surface participation ratio (SPR) classes and two target external coupling rates, with target frequencies between 5.0 and 7.2 GHz. These specifications are implemented by varying the CPW dimensions and the feedline coupling geometry. At the coupling point, the edge-to-edge spacing between the feedline gap and the resonator end gap is fixed at 5 μm for all resonators, while the coupling strength is tuned by varying the resonator end-gap length.

On chip label	SPR	Q_{ext} target	Freq target (GHz)	Track (μm)	Gap (μm)	Coupling gap length (μm)
1	low	low	5	40	24	50
4	low	high	5.2	40	24	110
7	mid	low	5.4	20	12	20
10	mid	high	5.6	20	12	70
2	high	low	5.8	10	6	2
5	high	high	6	10	6	35
8	low	low	6.2	40	24	65
11	low	high	6.4	40	24	125
3	mid	low	6.6	20	12	30
6	mid	high	6.8	20	12	75
9	high	low	7	10	6	2.5
12	high	high	7.2	10	6	45

5.5.2. DEVICE FABRICATION

This section provides additional details of the fabrication process used for the tantalum CPW resonators presented in the main text. General fabrication procedures are described in Section 2.2. The devices were fabricated on 4-inch, double-side polished (DSP), c-plane sapphire (Al_2O_3) wafers with a thickness of 430 μm . A 100 nm-thick α -phase tantalum film was deposited by STAR Cryoelectronics. Prior to fabrication, the wafer was diced into 10 mm $\times$ 10 mm chips.

The fabrication process begins with the removal of dicing resist by soaking the chips in acetone, followed by rinsing in isopropanol (IPA) and a brief oxygen plasma descum to remove residual organic contamination. Electron-beam lithography was used to define the resonator patterns. AR-P 6200.18 electron-beam resist was spin-coated at 4000 rpm and baked at 160 $^{\circ}\text{C}$ for 2 minutes, yielding a resist thickness of approximately 800 nm. A conductive layer of Electra 92 was then spin-coated at 4000 rpm and baked at 90 $^{\circ}\text{C}$ for 2 minutes to mitigate charging during electron-beam exposure on the insulating sapphire substrate.

The resonator patterns, shown in Fig. 5.1(a), were written using a Raith EBPg-5200 electron-beam lithography system. After exposure, the Electra 92 layer was removed by immersion in deionized water with mild sonication for 2 minutes. The resist was developed in pentyl acetate for 60 seconds, followed by 5 seconds in o-xylene and 60 seconds in IPA, after which the chip was dried with a nitrogen gun. A post-exposure bake at 130 $^{\circ}\text{C}$ for 2 minutes was performed to improve the etch resistance of the resist.

Pattern transfer was carried out using reactive-ion etching (RIE) in a Sentech Etchlab 200 system. After pumping to base pressure, an oxygen plasma descum (30 s at 50 sccm O_2 , 100 W RF power) was performed to remove residual resist. The tantalum film was then etched using a CF_4 -based process (20 sccm CF_4 , 10 μ bar chamber pressure, 50 W RF power). Endpoint detection was performed on a large exposed tantalum region near the center of the chip to determine when the film had been fully removed, after which the etch was continued for an additional 10 seconds to ensure complete removal of the film across the chip. Under these conditions, the etch time for a 100 nm-thick tantalum film was approximately 120 seconds. Following the etch, an additional in-situ oxygen plasma descum (30 s at 50 sccm O_2 , 100 W) was performed prior to venting the chamber.

After etching, the remaining resist was stripped in a hot dimethylformamide (DMF) bath maintained at 88 °C for at least 1 hour. The chips were then sonicated in DMF for 5 minutes at high power, followed by sonication in IPA for an additional 5 minutes. A final oxygen plasma descum was performed to remove residual organic contamination. This sequence defines the reference fabrication flow used for all devices in this chapter.

POST-FABRICATION CLEANING PROCEDURES

For the Piranha clean, chips were immersed for 20 minutes in a 2:1 mixture of $H_2SO_4:H_2O_2$ to remove organic residues. The chips were then rinsed sequentially in three separate baths of water and dried with nitrogen. For the buffered oxide etch (BOE), chips were immersed for 20 minutes in a 1:7 buffered HF solution to remove native surface oxides. After etching, the chips were triple-rinsed in water and dried with nitrogen.

After fabrication and cleaning, the chips were mounted on printed circuit boards (PCBs), glued in place using GE varnish, and wirebonded for microwave measurements. Multiple wirebonds were placed between the chip ground plane and the PCB ground to ensure robust grounding. In addition, on-chip wirebonds were placed across the feedline and resonators to suppress slotline modes and maintain a well-defined ground reference. A photograph of a mounted device is shown in Fig. 5.6.

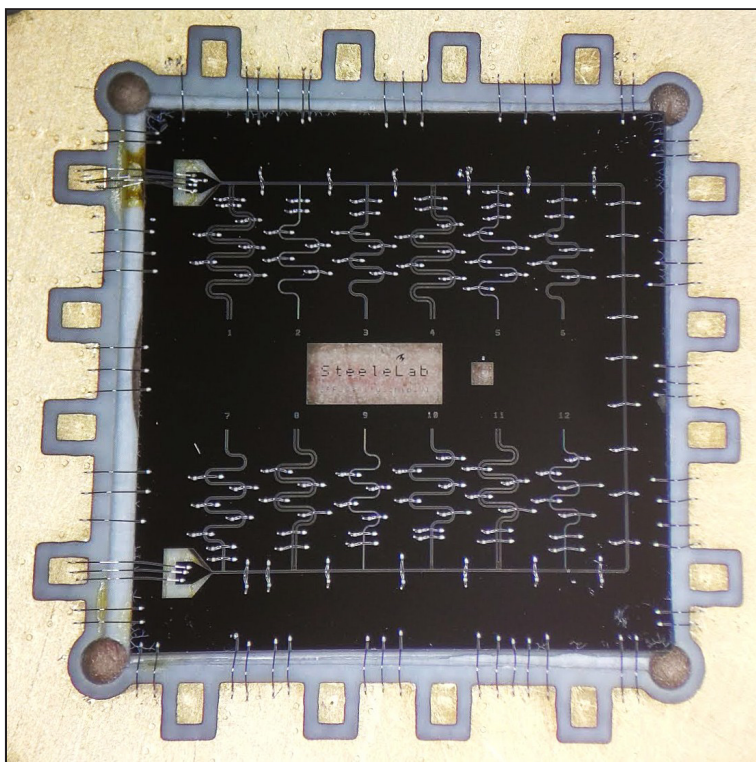


Figure 5.6: Photograph of a fabricated 10 mm $\times$ 10 mm tantalum resonator chip mounted on a printed circuit board (PCB) for microwave characterization. The chip is glued onto the PCB using GE varnish and electrically connected via aluminum wirebonds.

5.5.3. MEASUREMENT SETUP

The experiments were performed with the devices mounted on the base plate of a Bluefors dilution refrigerator operating at a base temperature of approximately 25 mK. A schematic of the cryogenic wiring is shown in Fig. 5.10. The signal input line has 53 dB of cryogenic attenuation, along with additional room-temperature attenuation that was adjusted during the experiments. The twelve linear resonators are coupled to a central feedline in a notch-type geometry, which connects to separate input and output lines. After passing through the chip, the transmitted signal is routed through a Low-pass filter (ZLSS-11G-S+) and a dual junction isolator (LNF-ISISC4_8a) before being amplified by a high electron mobility transistor (HEMT) amplifier (LNF-LNC4_8C) at the 4k stage, followed by a room temperature amplifier (Narda-MITEQ).

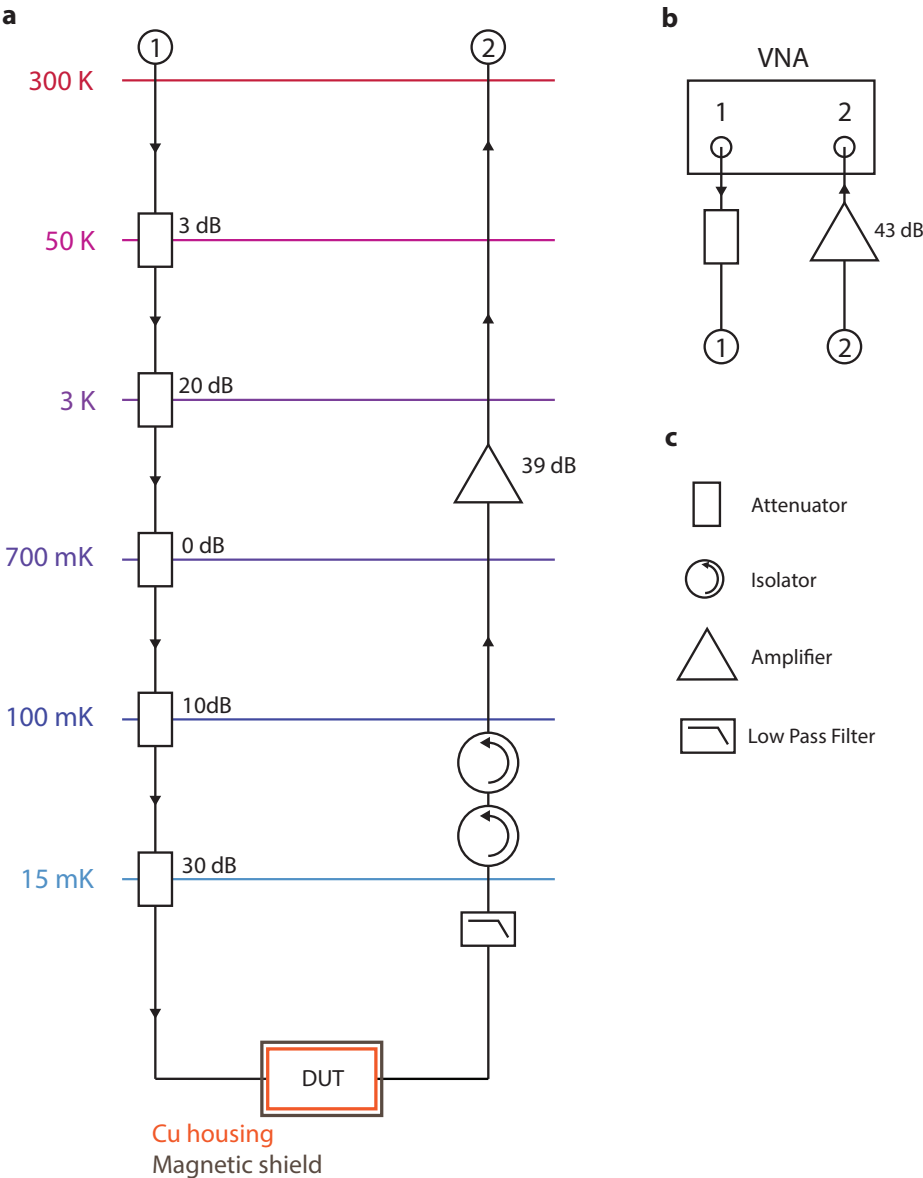


Figure 5.7: Schematic of the measurement setup. (a) Wiring diagram inside the dilution refrigerator. **(b)** Room temperature setup for VNA measurements. **(c)** Overview of components.

5.5.4. BACKGROUND CORRECTION OF COMPLEX TRANSMISSION DATA

To remove slow frequency-dependent amplitude and phase variations originating from the microwave measurement chain, we applied a complex background correction to the raw transmission data S_{21} . To access a larger power range than available from the VNA output, the measurements were performed in multiple sets with different room-temperature attenuation settings. The background of the transmission data was therefore corrected independently for each attenuation configuration.

For a given attenuation configuration, the complex transmission data were first averaged over all applied input powers to improve the signal-to-noise ratio of the background estimate. Two frequency windows far detuned from the resonance were then selected at the low- and high-frequency edges of the sweep, which was centered around the resonance frequency. These off-resonant regions were assumed to represent the background response of the measurement chain.

The background transmission was modeled by fitting the real and imaginary parts separately with first-order polynomials, yielding a complex-valued background model given by

$$S_{21}^{\text{bg}}(\omega) = (a_{\text{re}} \omega + b_{\text{re}}) + i (a_{\text{im}} \omega + b_{\text{im}}). \quad (5.2)$$

The measured transmission data were then normalized by this fitted background according to

$$S_{21}^{\text{corr}}(\omega, P) = \frac{S_{21}(\omega, P)}{S_{21}^{\text{bg}}(\omega)}. \quad (5.3)$$

This procedure removed slow frequency-dependent amplitude and phase variations while preserving the resonant response of the device. Applying the correction separately for each attenuation configuration ensured that attenuation-dependent changes in the measurement chain were removed. All transmission data presented in this chapter were background-corrected in this manner prior to fitting.

5.5.5. INPUT POWER CALIBRATION

To estimate the power at the device input, we calibrate the attenuation of the input lines and include the specified attenuation of additional cryogenic attenuators in the setup. The line attenuation is measured at a base temperature of $T = 25$ mK by connecting two input lines at the mixing chamber stage with a short coaxial cable and measuring the round-trip transmission. Figure 5.8 shows the resulting transmission data for all input lines in the dilution refrigerator. A polynomial fit to the measured round-trip transmission is used to obtain the average frequency-dependent attenuation, and the attenuation of a single line is estimated by dividing the round-trip value by two. This procedure captures the loss of the full signal path, including room-temperature cables and connectors.

The estimated power at the device input is then obtained by subtracting the estimated total attenuation from the source power,

$$P_{\text{dev}}(\omega) \approx P_{\text{src}}(\omega) - A_{\text{line}}(\omega) - A_{\text{att}}, \quad (5.4)$$

where $P_{\text{src}}(\omega)$ is the generator output power, $A_{\text{line}}(\omega)$ is the measured frequency-dependent attenuation of a single input line, and A_{att} is the fixed attenuation from additional attenuators.

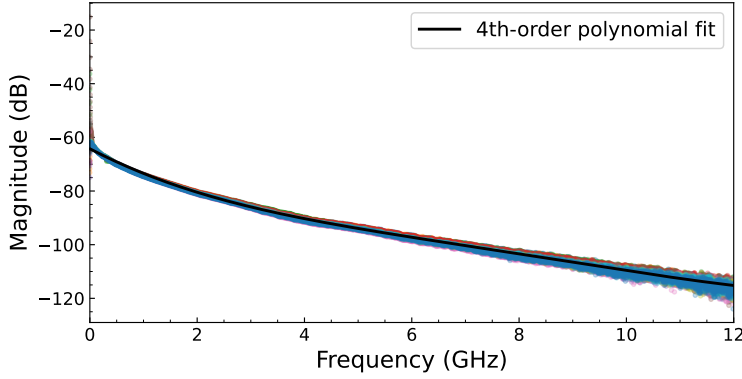


Figure 5.8: Estimation of the frequency-dependent input-line attenuation. Round-trip transmission of two connected input lines measured at $T = 25$ mK. The data show the frequency-dependent attenuation obtained for measurements of all input lines, together with a polynomial fit used to extract the average attenuation. The attenuation of a single line is estimated by dividing the fitted round-trip attenuation by two.

5

5.5.6. PHOTON NUMBER CALIBRATION

The average photon number in the resonator was estimated from the applied microwave power using the steady-state response of a driven hanger-type resonator. For resonant driving, the average photon number is given by [7, 32]

$$\langle n \rangle = \frac{2}{\hbar \omega_0^2} \frac{Z_0}{Z_r} \frac{Q_{\text{tot}}^2}{Q_c} P, \quad (5.5)$$

where ω_0 is the resonance frequency, Z_0 is the characteristic impedance of the feedline, Z_r is the characteristic impedance of the resonator, P is the microwave power at the device input, and Q_{tot} and Q_c are the total and coupling quality factors, respectively. Both impedances were taken to be $Z_0 = Z_r = 50 \, \Omega$, consistent with the CPW design and measurement environment. This calibration enabled the extracted resonance frequency and quality factors to be compared as functions of photon number. We note that this calibration provides an estimate of the average photon number and should be interpreted as an upper bound, because uncertainties in the total line attenuation and additional microwave losses between room temperature and the device are difficult to quantify precisely.

5.5.7. SINGLE-PHOTON Q_i EXCEEDING ONE MILLION

Fig. 5.9 shows the power dependence of one of the low surface participation ratio resonator for which an internal quality factor exceeding 10^6 was obtained, as shown in Fig. 5.2 in the main text. The extracted resonance frequencies and quality factors are plotted as a function of the average photon number. In the single-photon regime, the internal quality factor saturated above 10^6 as the photon number was reduced. The step-like change observed in both the resonance frequency and internal quality factor coincided with a change in the room-temperature attenuation and was attributed to changes in impedance matching and interference conditions in the measurement chain.

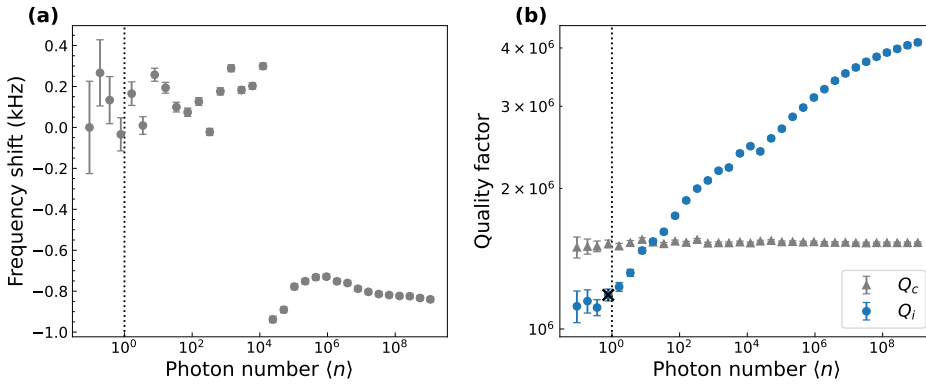


Figure 5.9: Extracted resonator properties as a function of photon number for a high- Q_i resonator. (a) Resonance frequency shift Δf as a function of the average photon number $\langle n \rangle$ for the low surface participation ratio resonator shown in Fig. 5.2, fabricated using the Piranha + BOE cleaning process. The frequency shift is shown relative to the resonance frequency $\omega_0/2\pi = 5.169$ GHz. **(b)** Internal (Q_i) and external (Q_c) quality factors extracted from fits to the complex transmission response as a function of photon number. The black cross marks the single-photon operating point corresponding to the fit shown in Fig. 5.2.

5.5.8. FULL DATA OVERVIEW

Figure 5.10 provides a complete overview of the measured resonance frequencies and quality factors for all resonators characterized in this work. Each resonator index refers to a specific resonator geometry, and shows data for devices subjected to different post-fabrication cleaning procedures. The resonators are ordered by increasing resonance frequency to facilitate comparison across the full frequency range.

For each resonator, the extracted quality factors are shown over the full range of measured photon numbers. Black crosses mark the single-photon internal quality factors obtained for under-coupled resonators, which form the basis of the cleaning-process comparison presented in the main text. The full data set illustrates the consistency of the surface-participation dependence and the reproducible improvement obtained with additional cleaning steps.

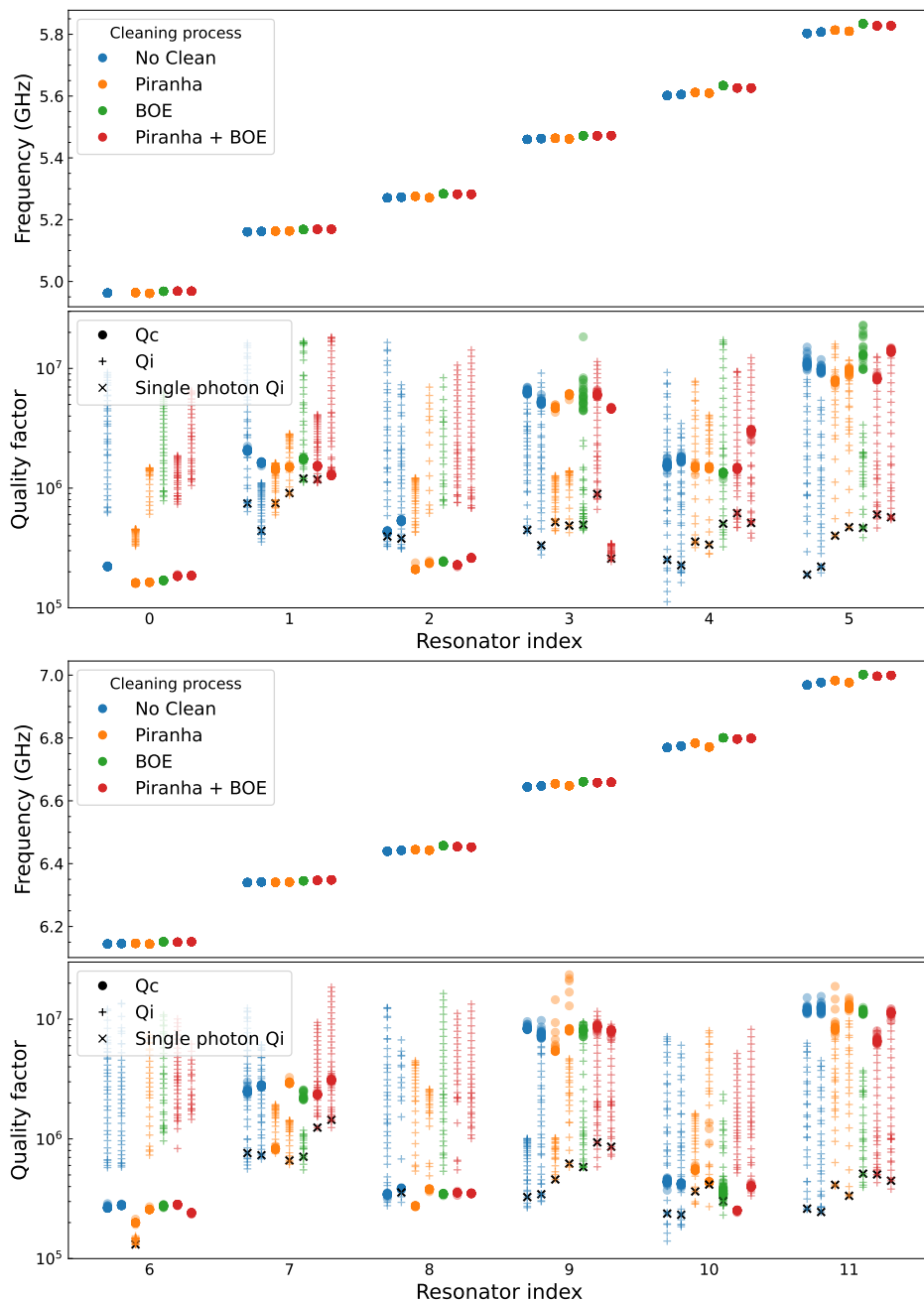


Figure 5.10: Overview of all measured resonator frequencies and quality factors. Measured resonance frequencies and quality factors are shown versus resonator index, ordered by increasing resonance frequency. The data includes all measured input powers. Black crosses mark single-photon internal quality factors extracted for the under-coupled resonators and used for the comparison presented in the main text.

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6

CONCLUSION AND OUTLOOK

6.1. CONCLUSION

In this thesis, we explored how superconducting microwave circuits can be improved by engineering interactions, implementing parametric amplification, and by reducing microwave loss. After an introduction of the used nanofabrication methods, we addressed several of the important practical details that go into fabricating superconducting circuits in the cleanroom.

In Chapter 3, we demonstrated how to engineer an intrinsic longitudinal interaction between a superconducting qubit and a linear microwave resonator. In contrast to effective longitudinal coupling schemes that rely on external drives, the interaction is directly encoded in the circuit design. Using a sideband drive, we demonstrated a dynamically controllable effective Jaynes-Cummings interaction in the strong-coupling limit with a high on-off ratio. By using the qubits ac-Stark shift as an absolute photon-number calibration, we extracted a longitudinal coupling rate that exceeded the largest dissipation rate in the system by approximately two orders of magnitude, which placed the system in the single-photon strong-coupling regime.

Chapter 4 presented the design, wafer-scale fabrication, and experimental characterization of a SNAIL-based parametric amplifier. We established its flux-tunable linear response and demonstrated stable parametric gain in the pumped regime, together with a characterization of the gain-bandwidth trade-off and the dynamic range. By combining details of the device design with the fabrication process and quantitative measurements, this chapter demonstrates how to obtain a parametric amplifier that can be directly integrated into microwave measurement chains for high-fidelity qubit readout and optomechanical experiments.

Finally, Chapter 5 addressed how to reduce dissipation in superconducting resonators. Using resonators fabricated from a single tantalum film, we compared single-photon internal quality factors across many devices as a function of resonator geometry and post-fabrication cleaning processes. We observed that additional post-fabrication cleaning produces a clear and reproducible improvement in the single-photon internal quality factor. In particular, a combined Piranha clean followed by buffered oxide etching yields the largest improvement, enabling reproducible single-photon internal quality factors exceeding one million for low surface participation ratio resonators. These results reinforce that surface-related dielectric loss is a major contributor to dissipation in superconducting circuits, and that targeted surface treatments can provide noticeable improvements in device quality factors without requiring changes to the underlying film.

6.2. OUTLOOK

The results of this thesis suggest several directions for future work.

The demonstrated intrinsic longitudinal coupling provides a possible route to coupling schemes that differ from conventional transverse coupling. A possible next step is to extend the architecture to multi-qubit devices, in which several qubits are coupled to a common linear resonator. In such a setting, frequency-selective sideband drives could be used to activate interactions only for chosen qubit-resonator sidebands, making it possible to selectively mediate qubit-qubit interactions. This approach may offer a path to on-demand couplings with reduced residual (static) ZZ interactions [1].

Another promising direction is to couple optomechanical systems to the linear resonator that mediates the longitudinal interaction. One of the key advantages of longitudinal coupling is its compatibility with large intracavity photon numbers. This is particularly relevant for optomechanics, where many protocols require strong coherent drives to enhance the radiation-pressure coupling, and to cool the mechanical mode to its ground state [2]. Coupling a qubit to an optomechanical cavity has been proposed for the preparation of mechanical quantum states [3–6]. However, in many of these implementations, the cavity inherits some Kerr non-linearity. This makes the strong drive regime difficult to reach without introducing unwanted nonlinear effects. Coupling a qubit longitudinally may offer a route to access these strong-drive regimes. This could provide a platform for the generation and control of macroscopic superposition states. Supporting the experimental investigation of decoherence mechanics related to gravity [7–9].

The demonstrated SNAIL parametric amplifier is directly suitable for integration into many different quantum measurement systems. Future experiments would benefit from an increase in the instantaneous bandwidth, and in an improved dynamic range (higher $P_{-1\text{dB}}$). A promising route to broader bandwidths is impedance engineering: by embedding the amplifier in an impedance-matching network, the gain profile can be flattened and extended beyond the bandwidth set by a single resonant mode [10, 11].

Finally, as superconducting circuits continue to improve and scale in complexity, material quality and surface-related two-level-system (TLS) loss are likely to become increasingly important [12, 13]. Within currently available material platforms, new strategies such as encapsulating superconducting films, have been proposed to suppress the formation of lossy surface oxides, which are now understood to be one of the dominant sources of TLS loss [14]. However, long-term progress will likely require advances beyond incremental process optimization. New material stacks and advanced interface-engineering techniques may be essential as device performance needs to improve further, and systems scale to larger architectures [15, 16]. Such advances would reduce physical error rates in superconducting quantum processors, easing the requirements for fault-tolerant quantum computation using quantum error correction [17, 18]. Overall, improvements in materials quality are likely to be a key enabling factor for the next generation of superconducting quantum hardware.

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LIST OF CONTRIBUTIONS

Journal Publications

4. B. Ann, S.-J. Choi, H.C. Park, S. Deve, **R.C. Dekker**, G.A. Steele, J. Ku, S.-B. Shim, and J. Suh, "Evidence for unexpectedly low quasiparticle generation rates across Josephson junctions of driven superconducting qubits", [arXiv:2511.22930](#), (2025).
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2. N. van Loo, G.P. Mazur, T. Dvir, G. Wang, **R.C. Dekker**, J.-Y. Wang, M. Lemang, C. Sfiligoj, A. Bordin, D. van Driel, G. Badawy, S. Gazibegovic, E.P.A.M. Bakkers and L.P. Kouwenhoven, "Electrostatic control of the proximity effect in the bulk of semiconductor superconductor hybrids", [Nature Communications](#) *14*, 3325, (2023).
1. G.P. Mazur, N. van Loo, J.-Y. Wang, T. Dvir, G. Wang, A. Khindanov, S. Korneychuk, F. Borsoi, **R.C. Dekker**, G. Badawy, P. Vinke, S. Gazibegovic, E.P.A.M. Bakkers, M. Quintero-Pérez, S. Heedt and L.P. Kouwenhoven, "Spin-mixing enhanced proximity effect in aluminum-based superconductor semiconductor hybrids", [Advanced Materials](#), *34*, 2202034, (2022).

Supervised student projects

- Matteo Arfini, "Microfabrication of a Superconducting Electromechanical Device Aimed at Enhancing the Mechanical Coherence of Drumhead Resonators", MEP Thesis, (Dec. 2022 - Sep. 2023).
- Pepijn Habing, "A numerical analysis of photon pressure coupling in cQED", BEP Thesis, (Apr. 2022 - Jul. 2022).
- Maël Demarets, "Driving-induced interactions in multi-nodes networks Towards reconfigurable networks of microwave qubits and photons", MEP Thesis, (Feb. 2022 - Aug. 2022).



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