Fast Qubit Control with a Quantum-Circuit Refrigerator

Timm Fabian Mörstedt
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Abstract

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Abstract
Superconducting circuits have emerged as powerful building blocks on the path toward a useful quantum computer. However, fast and accurate control over these circuits remains one of the key challenges. In particular, the fast initialization of superconducting qubits is a growing requirement in this era of constantly increasing qubit lifetimes.

In this thesis, we investigate different means of qubit control in the context of dissipation engineering. We use a quantum-circuit refrigerator (QCR), an on-chip microcooler based on one or two normal-metal–insulator–superconductor junctions, to create a tunable environment for superconducting circuits. We present and compare two different realizations of this device, the double-junction QCR directly coupled to a transmon qubit and the single-junction QCR coupled to the qubit via a superconducting resonator.

Beyond qubit reset, we explore other properties of the QCR, including the cooling and creation of exceptional points in superconducting resonators and the generation of thermal states in superconducting qubits. Through single-shot readout experiments, we gain insight into the quantum state of the qubit and its dynamics in response to different control signals. Combining the results of these experiments, we discuss the possible realization of a quantum heat engine using a QCR as a two-way tunable environment, extending the scope of applications toward the fundamental study of open quantum systems.

This thesis sheds light on the versatile world of quantum-circuit refrigeration and presents novel insights, experiments, and applications. At the intersection of circuit quantum electrodynamics and quantum thermodynamics, the QCR promises further possibilities for advancement and increased understanding of the behavior and control of superconducting quantum systems.

Keywords  quantum-circuit refrigerator, tunnel junction, photon-assisted tunneling, superconducting circuit, superconducting qubit, qubit reset


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It took exactly 30 minutes after sending my application email to receive a call back with an invitation to the Quantum Computing and Devices group at Aalto University. Of course, not everything can go this swiftly during a PhD project, but the spirit has stayed the same for the four and a half years that followed. The support that I have received during my PhD project has been overwhelming from all sides, and I want to express my gratitude toward all who made it happen.

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Espoo, June 11, 2024,

Timm Fabian Mörstedt
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This thesis consists of an overview and of the following publications which are referred to in the text by their Roman numerals.


Author’s Contribution

Publication I: “Rapid on-demand generation of thermal states in superconducting quantum circuits”

The author designed the sample circuit, carried out electromagnetic simulations on the sample circuit together with V.S, J.H., and J.R., fabricated the sample, carried out all measurements with assistance by W.S.T., A.G., and S.K., analyzed the data and wrote the manuscript with comments from all authors. Initial sample characterization was done by A.V., H.K., and M.T., alongside experimental discussions. W.S.T., M.R., V.V., G.C., and J.A. contributed through theoretical discussions. The work was supervised by M.M.

Publication II: “Recent Developments in Quantum-Circuit Refrigeration”

All authors contributed equally to the writing of the manuscript that was coordinated by A.V. and the author. M. M. supervised the work in general.

Publication III: “Initial experimental results on a superconducting-qubit reset based on photon-assisted quasiparticle tunneling”


Publication IV: “Single-junction quantum-circuit refrigerator”

V.V. carried out the theoretical investigation. V.V. and A.V. contributed to data visualization. All authors contributed equally to the writing of the manuscript. T.A.-N. and M.M. supervised the work conceived by M.M.

Publication V: “Exceptional-point-assisted entanglement, squeezing, and reset in a chain of three superconducting resonators”

W.S.T. carried out the theoretical investigation together with V.V. The author and S. K. contributed through discussions and circuit design. The manuscript was written by W.S.T. with comments from all authors. The work was supervised by M.M.

Publication VI: “Quantum-circuit refrigeration of a superconducting microwave resonator well below a single quantum”

A.V. conducted the experiments and analyzed the data. The author mainly developed the experimental setup. The author, W.S.T. and M.T. contributed to characterizing the sample which was designed and fabricated by the author. J.R. contributed to the simulations of the sample. M.S. supported in data analysis. The manuscript was written by A.V. and M.M. with comments from all authors. M.M. supervised the work.


W.S.T. conducted the experiments assisted by the author and analyzed the data. The author designed the sample circuit and fabricated the sample. The author, W.S.T., A.V., H.K., A.G., M.T., S.K.,
and A.S. contributed to the experimental setup. W.S.T. wrote the manuscript with comments from all authors. The work was supervised by M.M.
Further Publications by the Author


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Abbreviations

**cQED** circuit quantum electrodynamics  
**QCR** quantum-circuit refrigerator  
**NIS** normal-metal-insulator-superconductor  
**CPW** coplanar waveguide  
**BCS** Bardeen-Cooper-Schrieffer  
**DOS** density of states  
**RIE** reactive ion etching  
**EBL** electron beam lithography  
**AWG** arbitrary-waveform generator  
**TWPA** traveling-wave parametric amplifier  
**HEMT** high-electron-mobility transistor  
**IQ** in-phase-quadrature-phase  
**GMM** Gaussian mixture model  
**EP** exceptional point
1. Introduction

1.1 Motivation

With superconducting qubits emerging as powerful building blocks for quantum processors [1, 2, 3, 4], their fast and accurate control becomes one of the essential challenges on the track toward a useful quantum computer. One of the key operations of qubit control is the initialization of the qubit to its ground state [5]. Even though any qubit will naturally relax into its lowest-energy state over time by dissipating energy into its environment, this process can be very slow for high-coherence qubits as they are well isolated. And it gets worse: Superconducting qubits are constantly improving in terms of the lifetime of a quantum state, being able to hold excitations longer than ever before, currently reaching into the range of hundreds of microseconds [6] with millisecond lifetime on the horizon. This leaves us with the following conundrum: Even though it is desirable to achieve a high qubit lifetime to minimize computational errors, the wait time for natural decay of residual excitations between two computations grows proportionally and eventually slows down a quantum processor up to the point of impracticality. This problem increasingly demands a process that initializes the qubit quickly and accurately on demand.

Our criteria for a suitable reset protocol are 1) a short reset time that is significantly shorter than the time constant of natural decay, 2) high reset fidelity leading to minimal populations outside of the ground state, and 3) the criterion of unconditionality, requiring that a reset protocol can reach the ground state regardless of the initial state, even if it is outside the two-level computational subspace.

To this end, several different reset protocols have been proposed. Information-based methods [7], which require a non-demolition measurement of the qubit state to direct the qubit into its ground
state through a series of microwave pulses, are slow and susceptible to readout and gate errors. Unconditional microwave or resonator-assisted reset techniques such as $f_0-g_1$ [8] or parametric reset [9] can be significantly faster, but typically transfer the qubit excitation to another circuit element, for example to a resonator, and depend on its dissipative properties.

Dissipation engineering is a quite different approach to qubit initialization, and qubit control in general. Here, we aim to actively control the interaction of a quantum system with its environment. Based on circuit quantum electrodynamics (cQED) [10] of open quantum systems, the qubit can be placed in a dissipation-tunable environment that allows to preserve the quantum state in a low-dissipation setting and can rapidly remove excitations as a highly dissipative element.

In this thesis we use a quantum-circuit refrigerator (QCR), an electronic microcooler based on photon-assisted tunneling in one or two normal-metal-insulator-superconductor (NIS) junctions to create a tunable environment. These devices are well known for their use in thermometry [11, 12, 13, 14] and chip-level refrigerators [15, 16]. In the context of cQED, this standalone device has previously been coupled to superconducting resonators and proven its ability to provide cooling and tune the dissipation rate by several orders of magnitude [17, 18, 19].

For the purpose of qubit reset, the QCR can be utilized in two different ways: Direct reset, in which a simple dc voltage pulse to the QCR is sufficient to introduce dissipation in the entire coupled circuit, and hybrid reset [20, 21], where the operation of the QCR is assisting another reset protocol, for example a microwave-driven protocol by providing dissipation to another circuit element, for example a resonator. In both scenarios, the QCR can fulfill the criteria of fast reset, high reset fidelity, and unconditional reset of higher qubit states.

Beyond its application in qubit reset and refrigeration, the QCR lends itself to a more fundamental study of open quantum systems [18, 22, 23] and quantum thermodynamics. The former includes Lamb-shift experiments [24], whereas the latter enables studies on heat transport [25] and possibly the realization of a quantum heat engine [26].

1.2 Objective

Because this thesis is motivated by qubit reset, we first experimentally examine the direct reset protocol with a QCR directly coupled
to a superconducting qubit in Publication III. This proof-of-concept experiment not only shows the validity of this initialization method but also provides a baseline for the following experiments.

Based on this experiment and several studies on the QCR reviewed in Publication II, we introduce two significant changes to the system under study in Publications IV and V in order to improve its performance. Firstly, we simplify the device structure from a double-junction to a single-junction design and investigate its properties theoretically in the context of resonator reset and the generation of exceptional points. Secondly, we introduce a new coupling mechanism between a qubit and a QCR in order to improve the efficiency of qubit reset. This design revision makes use of a detuned resonator as a buffer between the qubit and QCR with the aim to preserve qubit coherence in the presence of the QCR.

These two changes trigger three different experimental studies: 1) The refrigeration of a resonator, which is probed through a transmon qubit, resolving its Fock state populations in Publication VI. This is not only the first experimental realization of a single-junction QCR, but also gives insight into the temperature of its electromagnetic environment. 2) The combination of a $f_0$-$g_1$ microwave drive-based reset protocol with QCR-based reset in Publication VII shows a different approach to qubit reset than Publication III. Here, the term $f_0$-$g_1$ denotes a transition from the second excited state of the qubit labeled by $f$ to the ground state labeled by $g$ under the excitation of a coupled resonator from zero to a single-photon Fock state. We also analyze the populations of higher qubit states outside of the computational subspace to gain insight into their dynamics. 3) The generation of thermal states in a transmon qubit in Publication I. This is somewhat adverse on the track toward qubit reset, but highly sought-after in the context of quantum heat transport [25], heat engines [27], and quantum thermodynamics. We combine the results from these experiments to discuss the realization of a quantum heat engine based on a QCR as a two-way temperature-tunable bath.

1.3 Thesis Structure

In Ch. 2, we introduce the theoretical concepts of superconducting circuits and quantum-circuit refrigeration, including an estimation of its key parameters. The fabrication process for resonators, qubits, and QCRs is described in Ch. 3, followed by a description of the experimental setup with a focus on single-shot readout. In Ch. 4, we discuss the double-junction QCR as in Publication II and its first experimental implementation in qubit reset in Publication III.
Chapter 5 serves as a platform to introduce the single-junction QCR with a discussion of theoretical results in Publications IV and V, followed by the results of experimental studies in Publications I, VI, and VII. Finally, Ch. 6 summarizes our findings and highlights potential future research and applications of this device.
2. Theoretical Background

This chapter introduces different physical phenomena and their theoretical description necessary for the understanding of dissipation control in superconducting quantum circuits. After a brief description of superconductivity, we delve into the main properties of coplanar waveguide (CPW) resonators and transmon qubits before focusing on the theoretical approach to quantum-circuit refrigeration. This includes a theoretical description of the QCR itself in Ch. 2.3.1, and, more importantly, its coupling to some of the elementary building blocks of circuit quantum electrodynamics in Ch. 2.3.2. More precisely, we first focus on the direct coupling of the QCR to a resonator and to a transmon qubit in Chs. 2.3.3 and 2.3.4, respectively. Subsequently, we discuss a more complex configuration where the QCR is indirectly coupled to the transmon through an auxiliary resonator in Ch. 2.3.5.

2.1 Superconductivity

Superconductivity, first discovered in 1911 by Heike Kamerlingh Onnes [28], is a thermodynamic phase in certain materials, characterized by the absence of electrical resistance and magnetic fields inside the superconducting matter [29, 30]. It only occurs below a critical temperature $T_c$ and a critical magnetic field $B_c$, where electrical resistance vanishes up to a critical current $I_c$. The phase transition under these conditions is described by the Bardeen–Cooper–Schrieffer (BCS) theory [31], which includes the formation of electron pairs, so-called Cooper pairs, due to electron–phonon coupling. The corresponding coupling energy $2\Delta$ defines the superconductor gap parameter $\Delta$ and affects the density of states (DOS) inside the superconductor. Even though the DOS vanishes completely within the energy range of $2\Delta$ according to BCS theory, experimental studies have found a small quasiparticle density inside the superconductor.
Theoretical Background

gap determined by the Dynes parameter $\gamma_D$ for physically realized superconducting materials [32]. The finite subgap DOS is caused in part by the lifetime broadening of quasiparticle excitations with the recombination time $\tau_r = \tau_0 \sqrt{\Delta/T} e^{\Delta/(k_B T)}$, with the Boltzmann constant $k_B$, and $\tau_0$ being a time-constant based on the electron–phonon coupling strength, and inelastic electron–electron scattering [33]. Further contributors to the subgap DOS are the inverse proximity effect [34] in superconductor–normal-metal hybrid structures, environment-assisted tunneling [35], two-electron Andreev tunneling [36], and external noise. Considering these effects in the Dynes parameter, the DOS at energy $\epsilon$ can be expressed as

$$n_S(\epsilon) = \left| \text{Re} \left[ \frac{\epsilon + i \gamma_D \Delta}{(\epsilon + i \gamma_D \Delta)^2 - \Delta^2} \right] \right|.$$  \hspace{1cm} (2.1)

The resonator and qubit circuits described in the following chapters are all based on superconductors. In this thesis we use Niobium (Nb, $T_c=9.26$ K, $B_c=0.82$ T) and Aluminum (Al, $T_c=1.2$ K, $B_c=0.01$ T) [37, 38, 39].

2.2 Qubit and Resonator Systems

2.2.1 CPW resonator

The inductor-capacitor resonator is one of the most important devices to study qubits and QCRs. In this thesis, it is used for qubit readout, as a QCR-resonator system, and as a coupler between the qubit and QCR. Due to its behavior as a harmonic oscillator, a resonator is one of the most basic components in quantum electrodynamics [40, 10, 41]. We will first look at its characteristic properties and resonator states, and then also at its interactions with other circuit components, most importantly qubits and QCRs.

The resonators are realized as quarter-wavelength coplanar waveguides (CPWs) of a given length $d$. As a harmonic oscillator, the Hamiltonian of a CPW resonator with capacitance $C_t$ can be written as

$$\mathcal{H} = \sum_{k=0}^{\infty} \frac{q_k^2}{2C_t} + \frac{C_t \omega_k^2 \phi_k^2}{2},$$  \hspace{1cm} (2.2)

with the charge $q_k$, flux $\phi_k$, and indices corresponding to the different eigenmodes of the system. The angular eigenfrequency corresponding to these conditions is

$$\omega_k = \frac{2k + 1}{\sqrt{L_t C_t}},$$  \hspace{1cm} (2.3)
with inductance \( L_k = 2L_r/(k^2\pi^2) \) and capacitance \( C_k = C_r \) of mode \( k \) [40]. Promoting the flux \( \phi_k \) and charge \( q_k \) to quantum operators, \( \hat{\phi}_k \) and \( \hat{q}_k \), and writing them in terms of annihilation and creation operators, \( \hat{a}_k \) and \( \hat{a}^\dagger_k \), we find the Hamiltonian operator

\[
\hat{\mathcal{H}} = \sum_{k=0}^\infty \hbar \omega_k \left( \hat{a}^\dagger_k \hat{a}_k + \frac{1}{2} \right). \tag{2.4}
\]

For the fundamental mode \( k = 0 \) we find the fundamental resonator frequency \( \omega_r = 1/(\sqrt{L_r C_r}) \) and [42]

\[
\hat{\mathcal{H}} = \hbar \omega_r \left( \hat{a}^\dagger \hat{a} + \frac{1}{2} \right). \tag{2.5}
\]

This single-mode version of the CPW resonator Hamiltonian is the one considered in the remainder of this thesis.

The most common states of the resonator here are Gibbs states and coherent states. Consequently, it is convenient to express the resonator states in the Fock basis using the occupation number \( n \) [43, 44, 45]. In the former case, the resonator state can be expressed as

\[
\hat{\rho} = \sum_{n=0}^\infty P_n |n\rangle \langle n|, \tag{2.6}
\]

with the probabilities of Fock states \( |n\rangle \) governed by the Bose–Einstein distribution \( P_n = \tilde{n}^n / (\tilde{n} + 1)^{n+1} \) and the average photon number \( \tilde{n} = 1/(e^{\hbar \omega_r/(k_B T)} - 1) \). On the other hand, the definition of a coherent state is

\[
|\alpha\rangle = e^{-|\alpha|^2} \sum_{n=0}^\infty \frac{\alpha^n}{\sqrt{n!}} |n\rangle, \tag{2.7}
\]

with Fock state probabilities following the Poisson distribution \( P_n = e^{-\tilde{n}} \tilde{n}^n / n! \) defined with the average photon number \( \tilde{n} = |\alpha|^2 \).

### 2.2.2 Transmon Qubit

Unlike the CPW resonator, the transmission line shunted plasma oscillation qubit, or transmon qubit, is an anharmonic system [1]. Derived from the charge qubit [46], it is based on a Cooper pair box, but sets itself apart by a large ratio \( E_J/E_C \gg 1 \) of the Josephson energy \( E_J \) and the charge energy \( E_C \), making the transmon robust against charge noise and dephasing, and therefore one of the most competitive designs among superconducting qubits at this point in time [6]. Its basic electronic circuit is shown in Fig. 2.1 and consists of a Josephson junction, which is a superconductor–insulator–superconductor junction, parallel to a shunt capacitor \( C_s \).

For a frequency-tunable transmon, the single Josephson junction
in Fig. 2.1 can be replaced by a superconducting quantum interference device (SQUID), a loop of two parallel Josephson junctions, which renders the qubit resonance responsive to the magnetic flux piercing this loop with periodicity $\Phi_0 = h/(2e)$ [47]. The flux-tunable frequency range is limited by the asymmetry of the two Josephson junctions $E_{J1}$ and $E_{J2}$ following

$$E_J(\Phi) = (E_{J1} + E_{J2}) \cos \left( \frac{\pi \Phi}{\Phi_0} \right) \sqrt{1 + \frac{(E_{J1} - E_{J2})^2}{(E_{J1} + E_{J2})^2} \tan^2 \left( \frac{\pi \Phi}{\Phi_0} \right)}.$$  \hspace{1cm} (2.8)

The Hamiltonian of a transmon has a similar form to that of a Cooper-pair box. We define the number and phase operators $\hat{n}$ and $\hat{\phi}$, the effective offset charge $n_g = (Q_r + C_g V_g)/(2e)$ with the gate voltage $V_g$, capacitance $C_g$, and environment-induced offset charge $Q_r$ to write

$$\hat{H} = 4E_C(\hat{n} - n_g)^2 - E_J \cos \hat{\phi}.$$  \hspace{1cm} (2.9)

By expanding the cosine in this equation up to the fourth order and assuming $E_C/E_J \ll 1$, we can calculate the eigenenergies of the transmon as

$$E_m = -E_J + \sqrt{8E_JE_C} \left( m + \frac{1}{2} \right) - \frac{E_C}{12} (6m^2 + 6m + 3),$$  \hspace{1cm} (2.10)

and identify the anharmonicity $\alpha = -E_C$ [1].

In this case, by introducing the annihilation and creation operators, $\hat{b}$ and $\hat{b}^\dagger$, the transmon Hamiltonian acquires the form

$$\hat{H} = \hbar \omega_q \hat{b}^\dagger \hat{b} + \hbar \frac{\alpha}{2} \hat{b}^\dagger \hat{b}^2,$$  \hspace{1cm} (2.11)

where $\omega_q = (\sqrt{8E_CE_J} - E_C)/\hbar$ is the transition frequency between its two lowest levels.

An ideal, perfectly isolated qubit can hold a quantum state infinitely long, but a physical qubit is always coupled to its environment, allowing exchange of energy. Over time, an excited qubit decays exponentially toward its lowest energy state $E_0$ on the energy decay timescale $T_1$. Similarly, the qubit drifts exponentially out of its original phase on the timescale $T_2$, typically referred to as the decoherence time. These decoherence parameters strongly depend on the environment and circuitry coupled to the qubit.

### 2.2.3 Dispersive Qubit Readout

In this work, the qubit is not used by itself, but always implemented in conjunction with a CPW resonator to read the qubit state. With a suitable coupling strength $g$ between qubit and resonator, the
Theoretical Background

Figure 2.1. The combined circuit of a QCR (blue), CPW resonator (light green), transmon qubit (purple), and readout resonator (dark green). The resonators couple to the qubit capacitively, whereas the QCR is in galvanic contact with the resonator.

qubit induces a state-dependent shift $\chi$ in the resonance frequency of the resonator [48]. It is important that the qubit and resonator frequencies $\omega_q$ and $\omega_{RO}$ are far detuned ($\delta = |\omega_r - \omega_q| \gg g$) so that the direct exchange of energy between the two components is suppressed. If these criteria are met and the two-level truncation of the transmon Hamiltonian is used, this interaction can be described by the Jaynes–Cummings Hamiltonian in the dispersive approximation,

$$\hat{H}_{\text{disp}} = \hbar \omega_{RO} \hat{a}^\dagger \hat{a} + \frac{\hbar}{2} \left[ \omega_q + 2 \chi \left( \hat{a}^\dagger \hat{a} + \frac{1}{2} \right) \right] \hat{\sigma}_z,$$

(2.12)

where $\hat{\sigma}_z$ is the $z$ component Pauli operator. The dispersive shift amounts to $\chi = \frac{g^2 \alpha}{\delta (\delta - \alpha)}$.

The qubit state can then be determined from the amplitude and phase of the field transmitted or reflected from the resonator [49]. It is important to point out that the dispersive readout is a quantum non-demolition measurement [50], because the measurement tone does not cause any excitation in the qubit [51]. However, the measurement unavoidably causes dephasing, and during readout, the qubit is subject to natural decay at the timescale of $T_1$.

2.3 Quantum-Circuit Refrigerator

In this section, we dive into the working mechanisms of the key device of this research work: the QCR. Beyond its properties as a standalone device, we discuss its effect when coupled to a resonator, qubit, or a more complex circuit.

Building on normal-metal-insulator-superconductor (NIS) general-
purpose refrigerators [11, 15, 52], the first realization of a QCR is a pair of NIS junctions in the form of an SINIS structure [17, 18]. Later studies from Publication IV onward use a simplified design based on a single NIS junction. For historic reasons we will focus on the double-junction device in this chapter, but the basic properties of the device are identical for both designs. Specific differences and calculations for the single-junction design are discussed in Sec. 5.1

2.3.1 Operation Principle

A bias voltage can be applied between the two superconducting electrodes and depending on this bias, quasiparticles can tunnel through the NIS junctions as collective excitations that carry the charge of a single electron. There are different mechanisms of tunneling, as depicted in Fig. 2.2 (a)–(c): At zero bias voltage, the electrons lack energy to tunnel through the junction and the tunneling rate is ideally zero. If a bias potential $eV$ smaller than the superconducting gap parameter $\Delta$ of the superconductor is applied, tunneling events are energetically suppressed, but possible under the absorption of a photonic excitation. This phenomenon of photon-assisted tunneling enables the QCR to absorb energy from its electromagnetic environment [18, 17]. On the other hand, if the bias potential is higher than $\Delta$, the emission of a photonic excitation is possible during the tunneling process, giving rise to applications such as a heater or a photon source [53]. In addition to the elastic tunneling processes considered in Fig. 2.2 (d) and Eq. (2.13) below, these photon-assisted tunneling processes can affect the measured current-voltage curve of the NIS junction mainly increasing the current by photon absorption in the subgap region and by photon emission above the gap. For voltages much higher than $\Delta/e$ this curve shows asymptotically ohmic behavior with a fixed tunneling resistance $R_T$.

The characteristic current of an NIS junction as shown in Fig. 2.2 d) as a function of voltage $V$ is described by

$$I(V, T_N) = \frac{1}{2eR_T} \int_{-\infty}^{\infty} dE n_S(E) [f(E - eV, T_N) - f(E + eV, T_N)],$$  \hspace{1cm} (2.13)

where we use the Fermi-distribution $f(E, T_N) = 1/[e^{E/(k_B T_N)} + 1]$. The temperature parameter $T_N$ refers to the normal-metal electronic temperature. The density of states is given by Eq. (2.1) [54, 55].

2.3.2 QCR Coupled to an Arbitrary Quantum Circuit

If the normal-conducting island or electrode of the QCR is capacitively attached to a superconducting quantum circuit through a coupling capacitor $C_c$, the voltage of the circuit mode couples capac-
**Figure 2.2.** Different tunneling processes through a single NIS junction. Green arrows denote tunneling processes that are possible, whereas red arrows indicate energetically suppressed tunneling events. Red waves represent photon emission and blue waves denote a photon being absorbed during the tunneling process. (a) Zero-bias case: If $V = 0$, no tunneling processes are possible. Even under the absorption of a photon, the electron energy is not sufficient to reach the band above the superconductor gap. (b) Photon absorption tunneling. If a low bias $0 < V < \Delta/e$ is applied between the electrodes, tunneling is only possible under the absorption of a photon. This scenario enables the cooling and reset of a coupled circuit. (c) High amplitude case. If $V \geq \Delta/e$, all three tunneling processes are possible, including photon emission. This can lead to heating and excitation of the coupled circuit. (d) Characteristic current-voltage curve of a QCR in the low-temperature (blue, 10 mK) and high-temperature (red, 500 mK) case according to Eq. (2.13) for the following parameters: $\Delta/e = 215 \, \mu$V, $\gamma_D = 5 \times 10^{-3}$, and $R_T = 30 \, k\Omega$. 
itively to the normal-conducting part of the QCR and hence provides an electrostatic coupling to the electron tunnelling in the QCR. Here, we encounter the first difference between the single-junction and the double-junction QCR. Although both couple to the circuit through the normal-conducting part, the double-junction QCR requires a separate connection to the ground plane after the second junction, but the normal-conducting electrode of the single-junction QCR can be attached directly to a circuit element, e.g., a quarter-wavelength resonator, that provides a drain to the ground plane.

For a QCR coupled to a quantum circuit, the total Hamiltonian can be split into contributions from the circuit $\hat{H}_0$ and the QCR [18, 23] as

$$\hat{H}_{\text{tot}} = \hat{H}_0 + \hat{H}_T + \hat{H}_N + \hat{H}_S.$$  \hspace{1cm} (2.14)

The QCR-induced terms are the tunneling Hamiltonian $\hat{H}_T$ and the Hamiltonians of the normal-metal and superconducting electrodes $\hat{H}_N$ and $\hat{H}_S$. The Hamiltonian of the coupled circuit $\hat{H}_0$ can be written as the sum of capacitive and inductive contributions $\hat{H}_C$ and $\hat{H}_\Phi$, which are already affected by the QCR through the coupling capacitor $C_C$. Without specifying the coupled circuit, the total charging energy can be considered as

$$\hat{H}_C = \frac{\hat{Q}_N^2}{2C_N} + \frac{(\hat{Q}_0 + \hat{\alpha}\hat{Q}_N)^2}{2\tilde{C}_0}$$

$$+ (\hat{Q}_0 + \hat{\alpha}\hat{Q}_N) \sum_{i=1}^M (\tilde{C}_V^{-1})_{0i}\hat{Q}_i + \frac{1}{2} \sum_{i,j=1}^M (\tilde{C}_M^{-1})_{ij}\hat{Q}_i\hat{Q}_j,$$ \hspace{1cm} (2.15)

where $\hat{\alpha} = C_C/C_N$ is the ratio of the coupling and normal-metal capacitances, $\hat{\Phi}_0$ and $\hat{Q}_0$ are the phase and charge operators of the circuit at the coupling point to the QCR, and $\hat{\Phi}_N$ and $\hat{Q}_N$ are the flux and charge operators at the normal-metal island. The matrix $\tilde{C}_M^{-1}$ and vector $\tilde{C}_V^{-1}$ depend on the coupling capacitance as well as the quantum circuit itself.

Quasiparticle tunneling processes through the NIS junction are accounted for in the tunneling Hamiltonian $\hat{H}_T$. First, we define the tunneling operator

$$\hat{\Theta} = \sum_{k,l,T} T_{lk} \hat{d}_{k\sigma}^\dagger \hat{c}_{l\sigma},$$ \hspace{1cm} (2.16)

to write the tunneling Hamiltonian as

$$\hat{H}_T = \hat{\Theta} e^{-i\tilde{\xi}(\hat{\alpha}\hat{\Phi}_0 + \hat{\Phi}_N)} + \hat{\Theta}^\dagger e^{i\tilde{\xi}(\hat{\alpha}\hat{\Phi}_0 + \hat{\Phi}_N)},$$ \hspace{1cm} (2.17)

where the tunneling probability is determined by the sum of the tunneling-matrix elements $\sum_{k,T} T_{lk}$, which is proportional to $1/R_T$. 


\( \hat{d}_{l,\sigma} \) is the creation operator of the normal-metal island and \( \hat{c}_{k\sigma} \) is the annihilation operator of the superconducting electrode. Of course, each tunneling process changes the accumulated charge within the normal metal electrode according to the term \( e^{\pm i \hat{\phi}_0 (\hat{\phi}_N + \hat{\phi}_0) / \hbar} \). Using [\( \hat{\phi}_N, \hat{Q}_N \] = \( i \hbar \), the charge transfer amounts to [56, 23, 18]:

\[
e^{\pm i \hat{\phi}_0 (\hat{\phi}_N + \hat{\phi}_0) / \hbar} \hat{Q}_N e^{\mp i \hat{\phi}_0 (\hat{\phi}_N + \hat{\phi}_0) / \hbar} = \hat{Q}_N \mp e. \tag{2.18}
\]

The last two contributions to \( \hat{H}_{\text{tot}} \) are the Hamiltonians of the normal-metal and superconducting electrodes \( \hat{H}_N \) and \( \hat{H}_S \), which include the microscopic electronic degree of freedom of the electrodes, and are offset by the bias energy \( eV \).

The effect of the QCR on a coupled quantum circuit is best assessed by studying the transition rates \( \Gamma_{mm'}(V) \) between different eigenstates of the circuit as well as the effective QCR temperature \( T(V) \). These functions are usually obtained by assuming \( \hat{H}_T \) to be a weak perturbation using Fermi’s golden rule. The transitions from eigenstate \( |m\rangle \) to \( |m'\rangle \) of the core circuit Hamiltonian are [18]

\[
\Gamma_{mm'}(V) = |M_{mm'}|^2 \frac{2R_K}{R_T} \sum_{\tau = \pm 1} F(\tau eV + \hbar \omega_{mm'} - E_N). \tag{2.19}
\]

The function \( F(E) \) calculates the normalized rate of single-electron tunneling events and \( M_{mm'} \) is the matrix element of the charge shift operator of the core circuit, \( e^{-i e \hat{\phi}_0 / \hbar} \). In Chs. 2.3.3 and 2.3.4, we calculate \( M_{mm'} \) for specific circuits. The definition of \( F(E) \) involves some parameters of the NIS junction, such as the density of states and the temperatures in the junction electrodes. Assuming identical temperature \( T_N \) in the superconducting and normal-metal electrodes, \( F(E) \) takes the form

\[
F(E) = \frac{1}{\hbar} \int d\epsilon n_S(\epsilon) \frac{f(\epsilon - E) - f(\epsilon)}{1 - e^{-E/(k_B T_N)}}. \tag{2.20}
\]

In this case, \( F(E) \) describes the forward tunneling rate, and \( F(-E) = e^{-E/(k_B T_N)} F(E) \) represents the backward tunneling process. Considering the two lowest eigenstates of the system, we can define the effective coupling strength

\[
\gamma(V) = \Gamma_{10} - \Gamma_{01}, \tag{2.21}
\]

and the effective temperature

\[
T(V) = \frac{\hbar \omega_{10}}{k_B} \left[ \ln \left( \frac{\Gamma_{10}}{\Gamma_{01}} \right) \right]^{-1} \tag{2.22}
\]

corresponding to the expected population in the two-level system. Notably, this temperature does not depend on the matrix element,
hence the transition frequency is the only parameter from the coupled quantum circuit. This allows us to investigate the on- and off-state temperatures without further system specifications. In the QCR off-state \( V = 0 \), we have the thermal equilibrium case \( T = T_N \). For the minimum temperature achievable with the QCR, we have to consider different scenarios, both within the subgap bias range, see Fig. 2.2 (b). The first scenario is the low normal-metal temperature case \( T_N \) below the crossover temperature \( T_{N}^{co} \). Here, we define the on-state voltage as

\[
eV_{\text{max}} = \Delta + \hbar \omega_{10} - e\delta V,
\]

including a voltage shift

\[
e\delta V = k_B T_N \left\{ \ln \left( \frac{\sqrt{\pi} k_B T_N}{\gamma_D \Delta} \right) + \frac{1}{2} \ln \left[ \ln \left( \frac{\sqrt{\pi} k_B T_N}{\gamma_D \Delta} \right) \right] \right\}.
\]

In the limit \( T_N \to 0 \), the minimum QCR-induced temperature is [23]

\[
T_{\text{min}}(T_N=0) = \frac{\hbar \omega_{10}}{k_B} \left[ \ln \left( \frac{2\hbar \omega_{10}}{\gamma_D \Delta \left( 1 - \frac{\hbar \omega_{10}}{\Delta} \right)} \right) \right]^{-1}.
\]

The second case is the thermal activation regime, in which \( T_N > T_{N}^{co} \) is high enough for the Fermi distribution and the superconductor DOS to overlap significantly, and enable thermal excitations as the dominant mechanism of quasiparticle tunneling. In this case, the temperature reaches a minimum of \( T_{\text{min}} \approx T_N/2 \) for the on-state voltage \( eV = \Delta - \hbar \omega_{10} \). Therefore, the crossover temperature is given as \( T_{N}^{co} = 2T_{\text{min}}(T_N=0) \) [57].

Finally, the last case is the above-gap scenario delineated in Fig. 2.2 (c), where \( eV > \Delta \). Here, tunneling by photon emission is dominant over that by photon absorption and the QCR-induced temperature increases above the off-state temperature as [18]

\[
T_{\text{ag}}(V) \approx \frac{eV}{2k_B} \left[ 1 - \frac{2\pi^2 \Delta^2 k_B^2 T_N^2}{3(eV)^4} \right].
\]

### 2.3.3 QCR Coupled to a Resonator

We consider a CPW resonator as our coupled circuit with the Hamiltonian from Eq. (2.5), as studied by Silveri et al. [18], and follow their procedure to calculate the QCR-induced transition rates \( \Gamma_{mm'}(V) \) in Eq. (2.19) by replacing \( \omega_{mm'} \) by \( \omega_r(m - m') \). For the circuit Hamiltonian, we assume coupling to the resonator through the normal-conductor electrode, and consider the charge eigenstates \( |q \rangle \) with integer \( q \) fulfilling \( \hat{Q}_N |q \rangle = eq |q \rangle \) as
\begin{equation}
\hat{H}_0 = \sum_{q = -\infty}^{\infty} \sum_{m = 0}^{\infty} |qm_q\rangle \langle qm_q| \left[ E_N q^2 + \hbar \omega_r \left( m + \frac{1}{2} \right) \right].
\end{equation}

The standard approach to the calculation of the matrix elements \langle E'|\hat{\Theta}|E \rangle is described by Ingold et al. [58]. Following this method, we use Eq. (2.17) and the junction electrode eigenstate \langle E \rangle to write the transition matrix element

\begin{equation}
\langle E', q' m'_q|\hat{\Theta}|E, q m_q\rangle = e^{i\frac{\pi}{2} \rho t} \langle E'|\hat{\Theta}|E \rangle \langle q' m'_q|e^{-i\frac{\pi}{2} \rho N}|q m_q\rangle \\
+ e^{-i\frac{\pi}{2} \rho t} \langle E'|\hat{\Theta}^\dagger|E \rangle \langle q' m'_q|e^{i\frac{\pi}{2} \rho N}|q m_q\rangle,
\end{equation}

with

\begin{equation}
\langle q' m'_q|e^{i\frac{\pi}{2} \rho \hat{\Phi}_N}|q m_q\rangle = \delta_{q',q\pm 1} \langle m'_q| q m_q\rangle.
\end{equation}

Starting from the charge-shifted resonator states \langle m_q\rangle = e^{-i\omega_q \hat{\Phi}} |m\rangle and the definition \langle m'_q| m_q\rangle = M_{mm'}^2, we obtain

\begin{equation}
M_{mm'}^2 = \left\{ \begin{array}{ll}
e^{-\rho (m-m') \frac{m'}{m}} \left[ L_{m'}^{(m-m')}(\rho) \right]^2, & \text{if } m \geq m' \\
e^{-\rho (m'-m) \frac{m}{m'}} \left[ L_m^{(m'-m)}(\rho) \right]^2, & \text{if } m < m', \end{array} \right.
\end{equation}

based on the generalized Laguerre polynomials \( L_{m'}^{(m-m')}(\rho) \)[59, 60, 61] and the interaction constant \rho = \frac{C^2_{\alpha}Z_i}{C^2_N R_K} with the von Klitzing constant \( R_K = \hbar/e^2 \).

If we look specifically at single-photon transitions, i.e., transitions between the resonator eigenstates \langle m \rangle and \langle m \pm 1 \rangle, the transition rates simplify considerably. Considering \( M_{mm-1}^2 = \rho m \) and \( M_{mm+1}^2 = \rho (m+1) \), we write them as a function of the mean thermal occupation number \( N_T \) as

\begin{equation}
\Gamma_{mm-1} = \gamma_T (N_T + 1)m, \quad (2.31)
\end{equation}

\begin{equation}
\Gamma_{mm+1} = \gamma_T N_T (m + 1). \quad (2.32)
\end{equation}

We can treat the QCR as a thermal bath of temperature \( T_T \) coupled to a resonator with mean occupation \( N_T = 1/\{\exp[(h\omega_r/(k_B T_T))] - 1\} \) with the coupling strength \( \gamma_T \). The latter can be calculated as

\begin{equation}
\gamma_T = \frac{\pi \gamma}{\omega_r} \sum_{\ell, \tau = \pm 1} \ell F(\tau eV + \ell h\omega_r - E_N),
\end{equation}

where \( \gamma \) is defined as the asymptotic coupling strength for \( eV/\Delta \gg 1 \) as

\begin{equation}
\tilde{\gamma} = \frac{2R_K}{\pi R_T} \omega_r g_T \tilde{\alpha}^2 \rho.
\end{equation}
The bath temperature can be calculated as

\[ T_T = \frac{\hbar \omega_r}{k_B} \left[ \ln \left( \frac{\sum_{\tau=\pm 1} F(\tau eV + \hbar \omega_r - E_N)}{\sum_{\tau=\pm 1} F(\tau eV - \hbar \omega_r - E_N)} \right) \right]^{-1}. \quad (2.35) \]

The single-photon transition rates conserve detailed balance \( \Gamma_{10}/\Gamma_{01} = e^{\hbar \omega_r/(k_B T_T)} \).

In a physical circuit, the resonator experiences further sources of dissipation apart from the QCR, for example through input or output lines or to other coupled circuits. With an external coupling strength \( \gamma_{\text{ext}} \) and a corresponding occupation \( N_{\text{ext}} \), the overall steady-state thermal occupation number of the resonator is summed up to

\[ N_T = \frac{\gamma_{\text{ext}} N_{\text{ext}} + \gamma_T(V) N_T(V)}{\gamma_{\text{ext}} + \gamma_T(V)}. \quad (2.36) \]

### 2.3.4 QCR Coupled to a Qubit

Let us examine the properties of a QCR capacitively coupled to a superconducting qubit as shown in Fig. 2.3. This system was theoretically studied by Hsu et al. [57] and we follow their results in this section. Many properties of the QCR are retained from the calculation above, but the circuit Hamiltonian is different. The Hamiltonian of the coupled circuit \( \hat{H}_0 \) includes the transmon circuit from Eq. (2.9), the charge energy of the QCR, and the coupling term between the qubit and the QCR.

Equations (2.22) and (2.25) demonstrated above that the thermal properties of the QCR only depend on the qubit frequency \( \omega_{10} \), independent of the architecture and other parameters of the qubit. For the transition rates, however, we consider the specific case of a transmon qubit and calculate the transition rates in the off and on-state of the QCR separately. Its characteristic energy ratio
$E_1/E_C$ enables us to approach the transmon as a weakly anharmonic oscillator.

In the off-state, we again have $T = T_N$ and $k_BT_N \ll \hbar \omega_{10}$ and the detailed balance $\Gamma_{10}^{\text{off}}/\Gamma_{01}^{\text{off}} = e^{-\hbar \omega_{10}/(k_BT)} \ll 1$. Using Eq. (2.19), we find

$$\Gamma_{10}^{\text{off}} = \bar{\gamma}\gamma_{D}\frac{\Delta}{\sqrt{\Delta^2 - (\hbar \omega_{10})^2}} \approx \bar{\gamma}\gamma_{D}. \quad (2.37)$$

Due to the finite Dynes parameter, the QCR-induced transition rate is finite even in the QCR off-state. In order to preserve the lifetime or an excited state, it is necessary to achieve low values of $\gamma_{D}$. Typical values for this parameter are in the range $10^{-5}$ to $10^{-3}$.

For the corresponding transition rate in the QCR on-state, we must again consider the low-temperature ($T_N < T_N^{\text{co}}$) and the thermal activation regime ($T_N > T_N^{\text{co}}$) separately. These regimes allow for qubit reset, because in the on-state, the transition rates $\Gamma_{10}$ can greatly exceed the natural energy decay time of a typical transmon. For thermal activation, we find the transition rate

$$\Gamma_{10}^{\text{on}} = 0.38 \times \frac{\bar{\gamma}\sqrt{k_BT_N\Delta}}{\hbar \omega_{10}}, \quad (2.38)$$

and the reset infidelity at a given reset time $t_r$ is

$$1 - F_r \approx e^{-\Gamma_{10}^{\text{on}}t_r} + \frac{\Gamma_{\uparrow}}{\Gamma_{10}^{\text{on}}}, \quad (2.39)$$

where $\Gamma_{\uparrow}$ is the total qubit excitation rate including the QCR and other sources. We can compare this relaxation rate and reset infidelity to the low-temperature case. For $T_N < T_N^{\text{co}}$, the transition rate changes to

$$\Gamma_{10}^{\text{on}} = \bar{\gamma} \frac{\sqrt{k_BT_N\Delta}}{2\hbar \omega_{10}} \sqrt{(eV_{\text{max}} + \hbar \omega_{10})^2 - \Delta^2}, \quad (2.40)$$

and we can substitute $\Gamma_{\uparrow} \approx \Gamma_{01}^{\text{on}}$ in Eq. (2.39).

A numerical study by Hsu et al. [23] comparing the low-temperature regime ($T_N = 10\text{ mK}, V = V_{\text{max}}$) and the thermal activation regime ($T_N = 100\text{ mK}, V = (\Delta - \hbar \omega_{10})/e$) has found a 14 times higher on-state relaxation rate and only five times higher excitation rate for the low-temperature case. Hence, the low-temperature operation of the QCR is advantageous for qubit reset, as it yields faster reset times and higher reset fidelities than the high-temperature case.

As a final measure of QCR performance, we write the QCR on/off ratio from Eqs.(2.37) and (2.40):

$$R_{\text{on/off}} = \frac{\Gamma_{10}^{\text{on}}}{\Gamma_{10}^{\text{off}}} \approx \sqrt{\frac{\Delta}{\hbar \omega_{10}^2\gamma_{D}}}. \quad (2.41)$$
2.3.5 QCR Coupled to a Resonator and a Qubit

Direct coupling of a QCR to a transmon qubit can already produce a fast qubit reset, but we also study a different circuit where the QCR is coupled to a CPW resonator, which is again coupled to a qubit. Even though this reduces the effective coupling between qubit and the QCR, it can be advantageous in the protection of $T_1$ in the off-state and in the application of different reset protocols. This corresponds to the full circuit in Fig. 2.1, which is the basis for the experimental results detailed in Ch. 5. It is key that the coupling resonator is far detuned from the qubit to reduce dissipation, similar to a Purcell filter.

The interaction with the coupled circuit is not fundamentally different from our calculations in the previous two chapters. Since the QCR is coupled to the resonator more strongly than to the qubit, it can cause high decay rates in the resonator as calculated in Sec. 2.3.3. The indirect coupling between the QCR and the qubit does not pose an impediment for qubit reset, since it can be further optimized with external control, for example, through parametric flux modulation [9] or microwave driving [20, 21]. The latter has been investigated in Sec. 5.4 through a QCR-assisted two-tone reset. By applying a voltage bias to the QCR while driving the $f_0\rightarrow g_1$ and $e\rightarrow f$ transitions in the resonator-transmon system, the performance of the reset can be considerably improved. This occurs because the two-tone drive pulses generate a transition from the second excited qubit state to the ground state under the emission of a photon into the highly dissipative resonator. Even though this transition can be driven with a short pulse below 100 ns [8], the protocol is ultimately limited by the decay time of the resonator.

As a final remark, the QCR can also affect the transition frequencies of the coupled circuits in the form of Lamb and ac Stark shifts as the QCR changes the electromagnetic environment of the circuit, renormalizing its transitions [23, 24]. The Lamb shift is caused by zero-point fluctuations of the environment, whereas the Stark shift originates from excitations in the environment and is therefore temperature-dependent. We do not calculate these shifts here since they are typically calibrated experimentally for different drive amplitudes.
3. Methods

3.1 Sample Fabrication

A major part of this thesis is the fabrication of accurate and reliable QCR devices, along with the other circuit components, most importantly SQUIDs. While SQUIDs have been optimized previously and generally yield good devices, the fabrication of NIS junctions is the major challenge on the path to successful qubit reset experiments.

The fabrication process can be divided into three parts: A patterned Nb layer that forms the majority of the circuit, two SIS junctions (Al/Al$_{ox}$/Al) per qubit SQUID and one or two NIS junctions (Cu/Al$_{ox}$/Al) for the QCR. Both junction types are deposited by shadow evaporation with a strategically designed overlap of layers from different deposition angles.

The niobium circuits, which include bond pads, waveguides, resonators, the qubit island, and couplers, as well as the ground plane, are fabricated by VTT with a well-established optical lithography and dry-etching process. A continuous 200 nm film is sputtered onto an HF-cleaned, intrinsic, high-resistivity 6" silicon wafer. An optical lithography mask patterned with the niobium circuit is used to expose optical AZ 5214E resist. Excess niobium outside of the desired pattern is removed with reactive ion etching (RIE), followed by the removal of the optical resist. After transfer from VTT, the wafers are covered in a protective layer of AZ 5214E resist and diced into smaller samples, each containing between four and nine $10 \times 10$ mm chips for individual fabrication batches. Afterwards, the resist is removed and the sample is cleaned in Acetone overnight.

The Josephson junctions, two for each SQUID, are fabricated using two-angle shadow evaporation and in-situ oxidation. First, a methyl methacrylate copolymer EL11/poly(methyl methacrylate) 950k A4 bilayer resist is applied to the sample by spin-coating, us-
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ing 1550 rpm for 60 s followed by 60 s of baking at 160 °C for each layer. Resist patterning is done by electron beam lithography (EBL) with a pattern as shown in Fig. 3.1. Since the exposed area for SQUIDs is relatively small, proximity error correction is not used here, but the dose is optimized at 1800 µC/cm² for development in a solution of methyl isobutyl ketone (MIBK) and isopropyl alcohol (IPA) with a ratio of 1:3 and a solution of 2-methoxyethanol and IPA (1:2) for 20 s each. This process was later changed to 1400 µC/cm² and development in MIBK:IPA (1:3) for 60 s and IPA for 60 s, as this allows for more accurate junction sizes due to longer development times and lower relative operator errors during development.

The aluminum electrodes are deposited using two-angle evaporation in a Dolan-bridge design. To remove any oxides on top of the niobium leads for a good aluminum-niobium interface, the surface is cleaned by 30 s of argon sputtering with an Ar pressure of 4.5×10⁻⁴ mbar and an acceleration voltage of 1 kV at 0° angle of incidence. The first aluminum layer is deposited at an angle of −20° to create the first electrode with a rate of 0.3 nm/s and a thickness of 30 nm. After 5 min of cooldown, the evaporation chamber is flooded with oxygen up to a pressure of 3.3 mbar to oxidize the surface of the bottom electrode. The oxidation pressure and time determine the final junction resistance for a given junction size. Different approaches from fast, high-pressure oxidation (1 min, 6.6 mbar) to a slow, low-pressure process (5 min, 2.0 mbar) were tried, with the latter providing a more uniform resistance distribution across several fabrication runs. However, direct resistance comparisons of different batches are difficult due to small variations in junction area from EBL and development. Another uncertainty is created by the termination of the oxidation process, which is given by the rate at which oxygen is pumped out of the chamber, which depends on its initial pressure. The effective oxidation time is therefore longer than the times given here.

The second aluminum layer is deposited at +20°, creating an offset to the first layer. Using the same crucible and evaporation rate, a thickness of 60 nm is required for the second layer to ensure good contact across the 200-nm sidewalls of the niobium structures at the given angle. To passivate the aluminum surfaces from natural oxidation after fabrication, another in-situ oxidation process is performed after deposition of the second electrode at a higher oxygen pressure of 11.5 mbar for 5 min. This creates a closed film of Al₂O₃, which prevents further surface degradation aiming to increase device lifetime. After evaporation, the sample is submerged in acetone for 12 hours to dissolve the resist and lift off excess aluminum. The device is assessed under an optical microscope before further processing.
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Figure 3.1. Fabrication and device issues. (a) Sample overview for a double-junction QCR capacitively coupled to a qubit. The transmon qubit (center) is coupled to the readout resonator (bottom), the SQUID (top), and to the drive line (left). The QCR finger capacitor on the right couples to the QCR. (b) SEM micrograph of an NIS junction immediately after fabrication. The offset features of the two-angle are visible with the bottom layer (Al) at the top of the image and the top layer (Cu) ca. 300 nm below. (c) Device as shown in (b), but after four weeks stored in atmospheric conditions. The aluminum electrode looks very similar to (b), but the copper electrodes have diffused and formed a non-continuous, granular structure. (d) Not only the junctions themselves, but the whole copper island is subject to decay. Additionally, copper seems to migrate along aluminum lines to the ground plane aluminum pad (top right). (e) Measurement of the resistance distribution at room temperature across 45 NIS test junctions on a single chip using a manual probe station. These junctions include a 3-nm aluminum layer underneath the copper electrode.
The first QCR devices of the double-junction design were fabricated in a similar way as Josephson junctions, replacing the top electrode with a 60-nm copper layer. For a QCR, we prefer a higher resistance per unit area to suppress the degradation of the Dynes parameter owing to Andreev currents. Accordingly, the oxidation parameters are adjusted to reach target resistances around 30 kΩ. As already found for Josephson junctions, longer oxidation times at low pressure are favored here as well, typically 2.4 mbar for 10 min.

This technique can generate functional QCR devices, but several issues have been found, most importantly diffusion and oxidation of the copper layer within a short timeframe after fabrication, combined with the relatively low yield of devices of approximately 30%. Fig. 3.1 compares an NIS junction immediately after fabrication with the same junction four weeks later, stored in the atmosphere. Diffusion of the copper layer has caused the formation of droplets and an incontinuous, non-conductive film as a result, rendering the device non-functional. Several measures can be taken to reduce this effect during storage of the device, e.g. vacuum storage, but cannot alleviate the effect completely. Combined with the device’s sensitivity to electrostatic discharge during processing, handling, and measuring, this causes a very low yield of functional devices in the cryostat.

To address these issues, two changes were made to the device. Firstly, using a single-junction QCR design as described in Publication IV reduces the necessary yield by a factor of two. Secondly, the introduction of a 3-nm aluminum adhesion layer between aluminum oxide and copper greatly improves the interface to the copper layer. Despite being an aluminum film, due to the reverse proximity effect caused by its low thickness, the adhesion remains a normal-conductor and does not change the properties and current-voltage characteristics of the device. Furthermore, in-situ deposition of these layers ensures clean interfaces, minimizing the risk of parasitic NIS or NIN junctions. This technique has been used to fabricate the devices for Publications I, VI, and VII.

We propose two further changes to the QCR device to minimize thermal excitations in the qubit caused by the QCR. A copper bath as shown in Fig. 3.2 (a), a large volume of copper which is connected to the normal-conducting electrode of the QCR, can increase the heat capacity and heat flow of the device and potentially reduce the normal-metal temperature at the junction. Even though this might reduce the heating capability of the QCR shown in Publication I, it can be a step towards low-noise qubits and lower qubit temperatures, especially in the QCR off-state. The second proposed change is the introduction of a quasiparticle barrier in the aluminum electrode.
Methods

Figure 3.2. False-color SEM micrographs of the latest device improvements. (a) Copper baths (blue) extending from the normal-conducting electrode of the NIS junction with an area of $300 \times 300 \mu m$ on each side and a thickness of 100 nm. The copper bath can be fabricated within the same evaporation process as the NIS junction. (b) NIS junction with quasiparticle barrier fabricated by three-angle shadow evaporation. The rightmost finger is the thin aluminum section (10 nm), followed by the thick aluminum layer (20 nm) in the middle and the 60-nm copper layer (brown) on the left.

to prevent quasiparticles from tunneling back through the junction. With the superconducting gap parameter of aluminum depending on film thickness, a thin (e.g. 10 nm) aluminum section of length $\approx 300$ nm directly at the junction combined with a thicker (30 nm) layer for the rest of the electrode can create such a quasiparticle barrier. However, this technique requires a three-angle evaporation process (Al 20 nm at $-20^\circ$, Al 10 nm at $0^\circ$, Al 3 nm and Cu 60 nm at $+20^\circ$), see Fig. 3.2 (b). In the other direction, quasiparticles are already prevented from penetrating the resonator by the barrier at the aluminum-niobium interface.

At the time of publication of this thesis, test devices with both improvements have been fabricated and analyzed by SEM and optical microscopy, but not yet measured at low temperatures.

3.2 Experimental Setup

To select the best chip out of a fabrication batch, all samples are pre-characterized with a probestation and a source-meter unit by measuring room temperature junction resistances of all SQUIDs and QCRs against the ground. Since QCRs are susceptible to electrostatic discharge damage, they are frequently checked again after installation into the cryostat. The QCRs are characterized at low temperatures through a current-voltage measurement with a source measure unit to determine the tunneling resistance and the Dynes
Methods

Figure 3.3. Measurement circuit in and outside the cryostat as used for the experiments in Publications I, VI, and VII. The input of the readout line consists of an arbitrary-waveform generator (AWG) and microwave source. After 60 dB of attenuation, the signal reaches the sample and the transmitted signal is amplified by a traveling-wave parametric amplifier (TWPA) and a high-electron-mobility transistor (HEMT) amplifier before reaching the PXI measurement unit. The qubit drive consists of a similar setup as the readout input, whereas the flux and QCR-lines are each combining separate rf and dc input lines with a bias-tee at the sample stage. For applying the QCR dc bias, a thermocoaxial cable is chosen over a superconducting wire for its quality as a low-pass filter. Fast QCR net-zero pulses are applied through the corresponding rf line. For details of the sample circuit see Fig. 5.3.
parameter.

All following measurements are carried out inside a dilution refrigerator at a base temperature of 35 mK. For Publications I, VI, and VII, a measurement setup as depicted in Fig. 3.3 is used for qubit and resonator measurements.

3.3 Single-Shot Readout

In all experimental studies included in this thesis, dispersive readout as described in Sec. 2.2.3 has been used to determine the state of the transmon qubit. For this purpose, we apply a 2-µs long microwave pulse with a frequency within the readout resonator linewidth to the input port of the readout line. The amplitude of this pulse is calibrated for maximum readout fidelity in each experiment, but is typically in the microvolt range. This pulse interacts with the readout resonator, and the signal transmitted through the readout line is amplified through a traveling-wave parametric amplifier at the sample stage and by low-noise amplifiers at the 50-K stage and at room temperature. The dispersive shift causes the readout resonator frequency to shift for different qubit states, creating a change in the amplitude and phase of the transmitted signal [10]. For analysis, one must differentiate between averaged readout, where the phase or amplitude of many measurements is averaged and treated as a single value, and single-shot readout, where each measurement is time-integrated, but not ensemble-averaged, allowing to analyze each measurement shot individually. While averaged readout is primarily used in Publications III and VI, single-shot readout is the primary mean of qubit state analysis in Publications I and VII. The main advantage of single-shot readout is the ability to analyze qubit state distributions and probabilities by observing point clouds in the in-phase–quadrature-phase (IQ) plane for repeated measurements. This is particularly important when measuring higher energy levels outside the two-level subspace or when measuring at elevated temperatures.

One disadvantage of single-shot readout is that data acquisition can be slow, making this method susceptible to drifts in the experimental parameters, e.g. flux voltage, during measurements of several hours. Especially large parameter sweeps as shown in Publications I and VII can take up to 24 h. Therefore, samples are first characterized using averaged readout and only specific experiments use single shots.

After the temporal integration of each readout trace in the time
window of 400–800 ns within the 2-µs readout pulse, the result of a single-shot measurement presents as several distinguishable point clouds in the IQ-plane. Besides integration time, the position of the clouds depends on readout frequency and amplitude, both of which are calibrated experimentally for optimal readout fidelity [41, 62].

There are several analysis methods available for determining the demarcation between different qubit states in the IQ plane. Regardless of the analysis method, measurement data from the prepared \(g, e, f, \) and \(h\)-state is recorded and combined, generating four two-dimensional Gaussian point clouds, which are used for calibration. The first analysis method is binning and manual fitting of four two-dimensional Gaussian functions to the calibration data [63]. The possible squeezing of qubit states requires asymmetric Gaussian distributions. Although this is possible, the large number of fit parameters can make this method unstable for strongly overlapping point clouds.

A generally preferred method to manual fitting are machine learning methods such as linear discriminant analysis, quadratic discriminant analysis, support vector machines, or the Gaussian mixture model (GMM), see Fig. 3.4. These methods do not require binning of data and all lead to reasonable demarcations for the single-shot data collected in this study. In the end, GMM was selected as the preferred method as it does not only provide demarcations but also supplies parameters of the four Gaussian distributions.

For a repeated single-shot measurement, knowing the position and covariance matrix for each Gaussian distribution, we can count the number of points inside of the elliptical \(1\sigma\)-boundary. The Mahalanobis distance \(m\sigma\) is defined as the distance of points of equal probability from the center of the distribution, and takes an elliptical shape. The ratio of points in and outside of \(m\sigma\) for an \(n\)-dimensional Gaussian distribution is described by the generalized regularized incomplete gamma function \(Q(n/2, 0, m^2/2)\). Any positive value of \(m\) can be used, as long as distributions do not overlap. Smaller values minimize false counts, but provide lower overall statistics. For example, the \(1\sigma\) ellipsoid of a two-dimensional Gaussian distribution has a ratio of \(Q(1, 0, 0.5)=39.34\%\) of the total points associated with this qubit state lie within its boundary [64]. Counting only points inside each of the four \(1\sigma\) ellipses provides the probability of the corresponding qubit state after normalization up to the fourth level. This method can be extended to higher qubit states, but fidelities may decrease.
Methods

Figure 3.4. Comparison of analysis methods for single-shot readout using the same dataset. (a) Linear Discriminant Analysis generates straight demarcation lines. Computation is fast, but false counts are possible, in this case, especially between $g$ and $e$. (b) Dataset of (a) analyzed with Quadratic Discriminant Analysis. The $g$-$e$ demarcation is improved compared to (a). (c) The Support Vector Machine allows even more degrees of freedom for the demarcation lines. (d) The Gaussian Mixture Model does not directly predict the qubit state and demarcations but predicts the parameters of Gaussian distributions based on machine learning. The ellipses show the $1\sigma$ boundaries of each qubit state.
4. Double-Junction QCR and Direct Coupling

The circuit design used in Publication III, a double-junction QCR coupled capacitively to a transmon qubit, presents the starting point for experimental studies in this dissertation. It is the first implementation of qubit reset with a QCR, and while it does have several shortcomings, it serves as a successful proof-of-concept experiment.

4.1 Circuit and Sample Design

Building on previous experiments on QCRs capacitively coupled to resonators [17, 53, 24], this study focuses on the fast initialization and cooling of a capacitively coupled qubit. Although the simplest way of capacitive coupling is to attach the QCR to an already existing capacitor between the qubit and the readout resonator, this fixes the QCR-qubit coupling capacitance. Our experiments in Publication III and the finite on/off ratio calculated in Sec. 2.3 indicate that the qubit lifetime $T_1$ is very sensitive to the coupling strength, and a strongly coupled QCR may affect the qubit even in the off-state. Therefore, a separate capacitor with a finger length that can be adjusted between different fabrication batches is attached to the normal-metal island of the QCR, see Figures 3.1 (a) and 4.1 (c).

In this reset protocol, the QCR is biased with a dc pulse. In order to reduce any high-frequency noise reaching the qubit through the QCR line, a Pd filter is added to the QCR line on the chip. The filter is fabricated by EBL and evaporation of a 30-nm Pd film. It is designed as wide fins (capacitors) connected by narrow leads (resistors) to form an RC filter circuit. The resistive part of the filter can slightly elevate the electron temperature by a few tens of millikelvins if pulsed at 1 MHz.
Figure 4.1. Experiment and sample for the double-junction QCR. (a) Cryostat and sample wiring. (b) Circuit diagram for the SINIS junctions (green) and transmon qubit (red). (c) SEM micrograph of the Pd filter (blue), QCR (green), and qubit (red). (d) High-magnification SEM micrograph of a single NIS junction with the normal-conducting (green) and superconducting (turquoise) electrodes. Figure adapted from Publication III (CC BY).
4.2 Qubit Reset and Design Shortcomings

As expected from the direct coupling architecture, the qubit reset experiments in Publication III show a very strong effect of the QCR on the qubit. Because of this, very short reset times can be achieved. Fig. 4.2 shows the excited state decay for different QCR pulse amplitudes between 0 and $eV/(2\Delta)=1.14$. The zero-bias measurement represents the natural decay of the transmon in the presence of a QCR. Even though typical transmon samples of this design without a QCR reach lifetimes between 10–30 μs, this measurement yields a significantly shorter decay constant of $T_1=1.76$ μs, suggesting increased dissipation also in the QCR off-state, in part due to additional dissipation from the QCR bias line.

The short-time behavior of the excited state probability follows an exponential decay for low pulse amplitudes, but deviates strongly from an exponential function for amplitudes near the superconducting gap, even including a temporary increase in the range of 5–50 ns. This effect shown in Fig. 4.2 (b) can be brought forth by charge accumulation in the normal conducting island. The time constant for this effect can be calculated with the subgap resistance near the gap voltage, $R = 1 \text{ M}\Omega$, and the normal-metal-island capacitance, $C_g = 24.4 \text{ fF}$, as $\tau = RC_g = 24.4$ ns, which is the same timescale on which this effect occurs in the experiment. For very short pulses on the nanosecond scale, the bias potential across each NIS junction can be different due to the nonlinear charging of the normal-metal island through the first NIS junction, leading to complex pulse distortion. Additionally, the time constant of the bias tee in the QCR input line can further affect the pulse shape. Beyond pulse distortion, amplitudes close to $eV/(2\Delta)=1$ can cause excitations in higher qubit levels due to strong coupling and physical proximity of the qubit and QCR. This effect has been observed in Sec. 5.5 and Publication I, but is not investigated experimentally in Publication III. As the sharp rise in excited probability between 5–50 ns for amplitudes near the superconducting gap contradicts high reset fidelity, we choose a medium amplitude case $eV/(2\Delta)=0.57$ as the most suitable for qubit reset, reaching $(3 \pm 1)\%$ of its starting probability within 80 ns. At this amplitude, a reasonable exponential fit yields a lifetime of $T_1^{\text{on}}=10\pm4$ ns, 170 times shorter than in the off-state. It is important to note that these probabilities are derived from a Rabi experiment and do not account for higher qubit excitations or the residual thermal population. The latter is theoretically estimated using the off-state transition rates in Eq. (2.37) to reach approximately 5% in the QCR off-state.

Using the parameters from this experiment, we can calculate
the theoretical on-off ratio using Eq. (2.41) as $R_{\text{on/off}} = 5000$ at low electron temperatures. To reach this value, improved filtering of the QCR input line as well as a physical separation of the QCR and qubit are necessary.
Figure 4.2. Qubit reset with a capacitively coupled double-junction QCR. (a) Time evolution of the excited state probability up to 300 ns for QCR pulse amplitudes between 0 and $eV/(2\Delta) = 1.14$. (b) The same measurement as in (a), but shown for very short reset pulses up to 50 ns. Within 6 ns, a reduction of the excited state probability below 10% of the starting value is visible for amplitudes higher than $eV/(2\Delta) = 0.8$. Comparison with fits for exponential decay (dashed lines) shows that the decay is highly non-exponential for this amplitude range and probabilities increase significantly between 6 ns and 50 ns. Figure adapted from Publication III (CC BY).
5. Single-Junction QCR and Indirect Coupling

5.1 Concept and Theoretical Considerations

The concept of a QCR based on a single NIS junction is first examined in Publication IV in a QCR-resonator system. This design change is motivated by simplicity and charge noise reduction. Additionally, in the fabrication of double-junction devices asymmetries between two junctions are unavoidable, including the Dynes parameter and tunneling resistance. This causes different behavior at the same bias voltage, which is eliminated in a single-junction QCR. In order to reach the same cooling power, the single junction must have halved resistance compared to a double-junction QCR.

Charge accumulation in the normal-conducting island is an effect arising in double-junction devices due to tunneling events in and out of the normal conducting island, effectively changing its charge state and nearby trapped charges [57]. This charge accumulation, which happens on the timescale of tens of nanoseconds, can cause a temporary voltage overshoot at one of the two junctions. In Sec. 2.3, we have neglected this effect by treating the normal-metal charging energy as sufficiently small by modeling the QCR as a bath with adjustable temperature and coupling strength. Even though this was in agreement with early QCR-resonator experiments [17, 19], this effect is suspected to contribute to the non-exponential decay during qubit reset in Fig. 4.2. This behavior does not occur in experiments with single-junction QCRs. With the single-junction design, the normal-metal electrode is directly connected to the superconducting input line or a superconducting grounded circuit element to prevent charge buildup. In this study, a quarter-wavelength resonator provides a direct channel to the ground potential for low frequencies.

The operation principle of a single-junction QCR is identical to
that of a double-junction device. The effective bias voltage per junction must be equal in the two cases, and hence the total pulse amplitude must be divided by a factor of two for the single-junction QCR for slow pulses. In conjunction with a quarter-wavelength CPW resonator as depicted in Fig. 5.1 (a), the QCR-induced dissipation rate and Lamb shift can be calculated with a semiclassical model for a resonator with low characteristic impedance $Z_r$, requiring $\sqrt{\pi Z_r/R_K} \ll 1$. Derived from Kirchhoff’s law in a circuit including the blue and light green parts of the circuit depicted in Fig. 2.1, we use a classical treatment of the resonator as

$$\omega_L - i\gamma = -\frac{i}{2} G_{\text{NIS}}(V, \omega_r) \omega_r Z_r$$

(5.1)

and the junction admittance $G_{\text{NIS}}(V, \omega_r)$ derived by Tucker et al. [65] to plot the voltage-dependent dissipation rate and coupling strength in Fig. 5.1 (b)-(e). As expected, the dissipation rate reaches its maximum at $eV = \Delta$ and the largest frequency shifts are observed at $eV = \Delta \pm \bar{\omega}_r$. A comparison with data from previous double-junction experiments shows no systematic deviation. This can be attributed to the dissipation and frequency shift being proportional to the sum of the conductances of the NIS junctions in series. However, this equivalence is only given for a symmetric SINIS junction, which still fulfills Eq. (5.1), if tunneling resistance and the characteristic impedance are adjusted accordingly. This does not hold true for an asymmetric junction due to different bias potentials at each junction.

A different calculation is required for the quantum limit $\sqrt{\pi Z_r/R_K} \gg 1$, where multi-photon processes are no longer negligible. In this case, we can employ the Born-Markov theory, and the frequency shift and dissipation rate can be extracted from the eigenvalues of the Liouvillian in the master equation. Unlike the previous calculation, the high-impedance case shows several resonances for both $\omega_L$ and $\gamma$, see Fig. 5.1 (f)-(i), the number of which increases with reduced temperature. The bias energy spacing of these resonances is $eV = \Delta \pm n\bar{\omega}_r$ with $n$ being the number of photons absorbed during the associated tunneling process.

The QCR architecture in Fig. 5.1 (a) can also be used for a theoretical study of exceptional points in a resonator chain as an example of a controlled open quantum system. This system consists of three capacitively coupled CPW resonators, see Fig. 5.2 (a). Two of the resonators are coupled to their respective input lines through a single-junction QCR to provide in-situ-tunable dissipation rates. Tuning the dissipation rate for each resonator can create open system degeneracies, also called exceptional points (EPs) [66, 67, 68, 69, 70]. These differ from Hermitian degeneracies by the coalescing eigenvectors and eigenvalues of the dynamical matrix describing the evolution of
Figure 5.1. Single-junction QCR. (a) Schematic of a single-junction QCR coupled to a quarter-wavelength CPW resonator. (b),(c) Coupling strength and frequency shift of a low-impedance (34.8 Ω) resonator calculated from Eq. (5.1) at $T=170$ mK (blue line) and experimental data from Silveri et al. [22]. (d),(e) The same calculation as in (b),(c), but at $T=0$ K. (f),(g) Coupling strength and frequency shift calculated for the high-impedance case ($Z_r=20$ kΩ, $R_T=640$ kΩ) for $T=170$ mK and (h),(i) $T=20$ mK. The different lines originate from different resonances of the system, and the intensity $f_j$ is based on the contribution of the resonance to the response function. Figure adapted from Publication IV (CC BY).
the open system. At a critical point, the system exhibits critical time-
dynamics, leading to polynomial solutions in time \cite{71,72}. Tuning
of a resonator or qubit system to an exceptional point constitutes
a form of a optimized dissipative reset since this is the condition,
under which the systems shows the fastest decay into the ground
state.

In the case of weakly coupled resonators and initial squeezing
in the third resonator, the center resonator can be identified as a
fast squeezing splitter \cite{73} and a fast distant entanglement \cite{74}
generator near a specific second-order EP. Different behavior can
be observed in the vicinity of a third-order EP, where an accelerated
decay of excited states can be observed, see Fig. 5.2 (b). These
examples showcase the potential versatility of the device upon the
detailed knowledge of its exceptional points.

The frequency $\omega_1$ and dissipation rate $\gamma_1$ of the first resonator
remain constant in the studied setup and changes in the other
two resonators are described with the complex-valued parameters
$\{\epsilon_k\}$ defined as with the frequency detuning $\delta_k = \sqrt{2}g\text{Im}(\epsilon_k)$ and the
dissipation rate $\gamma_k = \gamma_1 + 2\sqrt{2}g\text{Re}(\epsilon_k)$, where $g$ is the coupling strength
between the resonators.

Six third-order EPs can be found in the circuit described above. If
all resonator modes are degenerate, two occur at $(\epsilon_2, \epsilon_3) = (\pm 1, \pm 2)$,
resulting in $\gamma_2 = \gamma_1 \pm 2\sqrt{2}g$ and $\gamma_3 = \gamma_1 \pm 4\sqrt{2}g$. For non-degenerate
modes, four more EP-3’s appear with decay rates $\gamma_3 = \gamma_1$, $\gamma_2 = \gamma_1\pm 3\sqrt{3}g$, and frequency detunings of $\delta_2 = \pm g/2$ and $\delta_3 = \pm g$. To create
these non-resonant EP-3’s in an experimental scenario, frequency
tunability in the range $g/2\pi$ is required, so that we only focus on the
degenerate case. Because the QCRs can only increase dissipation,
we use $\gamma_2, \gamma_3 > \gamma_1$, requiring $\epsilon_k \geq 0$. Further EPs can occur, but
are limited to second-order degeneracies. In particular, the fast
quasi-stabilization of entanglement and squeezing is obtained at an
EP-2 with $(\epsilon_2, \epsilon_3) = (2, 0)$, yielding the decay rates $\gamma_2 = \gamma_1 + 4\sqrt{2}g$ and
$\gamma_3 = \gamma_1$. The quasi-stabilization of these quantities is dictated by the
smallness of $\gamma_1$ and $\gamma_3$.

Fast decay of non-classical states toward the ground state can
be realized with both EP-2s and EP-3s. As an example, after finite
squeezing in the third resonator, the initial state is generated during
a preparation time $\tau_s$ at the aforementioned EP-2, and reset is
performed over the reset time $\tau_r$. In Publication V, reset infidelities
have been calculated for $\tau_s = 0$ and $\tau_s = 8/g$. In a parameter sweep
of $\text{Re}(\epsilon_2)$ and $\text{Re}(\epsilon_3)$, regions of lowest reset infidelity tend to follow a
particular exceptional point branch passing through the EP-3. At the
EP-3, the reset infidelity reaches values below $10^{-5}$ within $\tau_r = 6/g$
independent of $\tau_s$ and the distribution of squeezing between the
Figure 5.2. (a) Circuit schematic of three coupled resonators for the creation of exceptional points. Each of the two resonators is coupled to a QCR for dissipation control. In addition, one resonator is frequency-tunable through a SQUID. (b) Reset infidelity for different preparation and reset times at a third-order exceptional point. Figure adapted from Publication V (CC BY).
different resonators, see Fig. 5.2 (b). This infidelity is further increased by a small initial squeezing parameter and a small ratio of $\gamma_1/g$.

### 5.2 Sample Design

For experimental studies on resonator and qubit control with a QCR, primarily for Publications I, VI, and VII, a sample design based on the full circuit depicted in Fig. 2.1 is developed in CAD, see Fig. 5.3 (a). The structures are drawn in KLayout using the KQ Circuit library, and electromagnetic simulations are conducted in Sonnet Suites, Ansys HFSS, and Qucs. For this purpose, the two-dimensional circuit geometry is exported from KLayout and imported into Sonnet Suites, where layer thicknesses and material parameters are specified. In these simulations, the circuit geometry, material layering, the SQUID resistance, the QCR tunneling resistance, and the QCR junction capacitance are input parameters to calculate the remaining values listed in Tab. 5.1. Fig. 5.3 (b) shows the simulated electric field in a subsection of the circuit including the qubit, reset resonator, and QCR at the reset resonator frequency. The circuit shown in Fig. 5.3 (b) already includes the Cu baths, which have a very small capacitance to the ground plane and therefore little effect on the simulation result.

In these simulations, the QCR resistance is set to either 15 kΩ to simulate the QCR on-state or 10 MΩ for the QCR off-state. At the qubit frequency of 5.2 GHz, the former results in a circuit quality factor of $10^4$, whereas the latter increases this value above $10^6$. Different NIS resistances do not affect the resonances in this simulation.

The circuit parameters and simulation results are shown in Tab. 5.1. The coupling capacitor dimensions are chosen depending on the resonance frequencies of the circuit elements for a given coupling strength. The coupling point of the qubit to the reset resonator is chosen close to the middle of the reset resonator since the ends of the resonator are already connected to the QCR and to the ground plane, respectively. The coupling point is chosen such that it does not create resonances that overlap with those of other circuit elements. Additionally, the potential qubit reset time is mostly determined by the reset resonator quality factor and coupling strength to the transmon.

The readout and reset resonator frequencies are chosen to be far apart, allowing for a qubit frequency to be placed in between them. With a flux-tunable qubit frequency it is possible to tune the qubit...
Figure 5.3. Sample design for indirect coupling of a QCR and qubit. (a) CAD-design for a circuit as in Fig. 2.1. The flux-tunable transmon qubit is coupled to a reset resonator with a coupling strength \( g_1 \) and to the readout resonator with a coupling strength \( g_2 \). This circuit is designed for a single-junction QCR (insert). The NIS junction attaches to the resonator end on one side and to the input line on the other. The direction of the junction is irrelevant for our simulations. (b) Simulation of the electromagnetic field in a circuit including the qubit, reset resonator, and QCR at the reset resonator frequency of 4.76 GHz using Ansys HFSS. The orange squares on either side of the QCR are Cu baths as proposed in Ch. 3.1. Simulation image courtesy of Jukka Räbinä/IQM.
Table 5.1. Simulated sample parameters.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Symbol</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Qubit frequency</td>
<td>$\omega_{ge}/2\pi$</td>
<td>5.20 GHz</td>
</tr>
<tr>
<td>Qubit anharmonicity</td>
<td>$\alpha/2\pi$</td>
<td>-273 MHz</td>
</tr>
<tr>
<td>Energy ratio</td>
<td>$E_J/E_C$</td>
<td>48.96</td>
</tr>
<tr>
<td>SQUID resistance</td>
<td>$R_{SQUID}$</td>
<td>10.5 kΩ</td>
</tr>
<tr>
<td>Critical current</td>
<td>$I_C$</td>
<td>27.28 nA</td>
</tr>
<tr>
<td>Reset resonator frequency</td>
<td>$\omega_1/2\pi$</td>
<td>4.67 GHz</td>
</tr>
<tr>
<td>Readout resonator frequency</td>
<td>$\omega_2/2\pi$</td>
<td>7.44 GHz</td>
</tr>
<tr>
<td>Reset coupling strength</td>
<td>$g_1/2\pi$</td>
<td>59.6 MHz</td>
</tr>
<tr>
<td>Reset coupling capacitance</td>
<td>$C_1$</td>
<td>4.8 fF</td>
</tr>
<tr>
<td>Readout coupling strength</td>
<td>$g_2/2\pi$</td>
<td>70.4 MHz</td>
</tr>
<tr>
<td>Readout coupling capacitance</td>
<td>$C_2$</td>
<td>3.3 fF</td>
</tr>
<tr>
<td>QCR tunneling resistance</td>
<td>$R_T$</td>
<td>15.0 kΩ</td>
</tr>
<tr>
<td>QCR junction capacitance</td>
<td>$C_{NIS}$</td>
<td>2.0 fF</td>
</tr>
</tbody>
</table>

In resonance with the reset resonator for maximum coupling. In addition to sample design considerations, the simulated frequency parameters are used as a starting point of the experimental characterization process of these samples.

Apart from the sample with parameters in Tab. 5.1, a similar circuit with a frequency-tunable reset resonator and greatly reduced reset resonator coupling capacitance has been devised for the experimental study of exceptional points. Even though this circuit differs from the three-resonator system described in Ch. 5.1, with the two resonators coupled via the qubit, exceptional points can still be generated in this system. Although typical qubit reset experiments can work within a range of different coupling strengths, for the operation at an exceptional point, the parameters must be accurately adjusted to fulfill $\delta = 0$ and $\gamma = 4g_1$.

### 5.3 Resonator Control and Cooling

In our first experimental realization of the single-junction QCR shown in Fig. 5.6 (a), the sample circuit described in Sec. 5.2 can be used to probe the reset resonator states with the transmon qubit. Specifically, we observe several spectral lines corresponding to the populations of different Fock states in the reset resonator, see Fig. 5.4 (a) and (c) as well as Fig. 5.5 (a) and (b). Additionally, we study the refrigeration of the resonator to a small fraction of a photon starting from a thermal state up to 1 K generated with
Figure 5.4. Resonator cooling as function of QCR bias voltage. (a) Measured phase of the transmitted signal through the readout line probed with a readout tone at the readout resonator frequency as a function of the QCD DC bias while applying a constant coherent drive of -130 dBm in resonance with the reset resonator. For $V_{QCR}=68 \mu V$, the inset on the right shows the corresponding Fock-state populations (bars) extracted through the relative intensities of the spectral lines. The orange dots represent a fit of the thermal distribution to the experimental data. (b) Photon number and decay rate extracted from (a) compared to a theoretical model using $C_{NIS}=0.54 fF$, $T_N=150$ mK and $Z_r=63.7 \Omega$ and Eq. (2.33). The crosses represent measurement results with an added QCR drive of 9.25 GHz. (c) As (a), but with additional heating by a strongly driven attenuator and without coherent drive of the reset resonator. (d) Mean photon number and effective temperature extracted from (c). The theory curve corresponds to Eq. (2.35). Figure adapted from Publication VI.

high-temperature artificial noise.

In addition to a DC bias already used in the experiments in Sec. 4.2, here the QCR can also be driven with an rf excitation using a measurement setup as described in Sec. 3.2 and Fig. 3.3. The results of this experiment are shown in Fig. 5.4. For this particular measurement, a continuous rf drive of 2.9 GHz was used for the QCR, but due to the broad bandwidth control of the QCR, the frequency choice is not critical and also allows for combined dc+rf drive signals.

In an experiment sweeping the qubit drive frequency and the reset resonator drive power, we can observe the spectral lines of different Fock states with their population corresponding to their phase amplitude. Assuming $\gamma_{QCR}^{off} < \gamma_{dr} < \gamma_{QCR}^{on}$, we can extract the mean resonator population and thus the resonator temperature, see Fig. 5.4.

A strong cooling effect can be seen in Fig. 5.4, reaching around
Figure 5.5. Artificially heated resonator cooled by a dc-biased QCR, measured with a similar technique as in Fig. 5.4. The qubit resonances are shown for a range of (a) artificial noise levels caused by an off-resonantly driven attenuator and (b) QCR dc bias voltages. (c) Mean photon number and effective temperature extracted from (b). Again, the theory curve is based on Eq. (2.35). Figure adapted from Publication VI.
100 mK at the gap voltage. Moreover, the increased decay rate in Fig. 5.4 (b) is in agreement with our calculations in Sec. 5.1 and Fig. 5.1. Further measurements with a 10 ns pulsed rf excitation instead of a dc bias voltage shows a comparable behavior to Fig. 5.4 (b) for drive strengths between $-80$ and $-60$ dBm.

To create thermal states in the reset resonator for subsequent cooling, we use an attenuator inside the cryostat at the sample stage driven with a strong, detuned signal to avoid interference with other circuit components. Due to black-body emissions from the attenuator, the photon number in the reset resonator is increased, which is equivalent to a rise in temperature, see Sec. 2.2.1. The corresponding population of the QCR environment can be calculated from Eq. (2.22). Under the continuous influence of band-limited artificial thermal noise with a power of $-105$ dBm, strong cooling from 1 K as low as 250 mK can be achieved, see Fig. 5.5. The measured photon number decreases more abruptly than predicted, possibly due to bias drifts during the experiment.

In these experiments we find a large operation bandwidth for the QCR bias signal, including dc, pulsed rf and continuous rf drives. As an additional step in this direction, this circuit could provide a platform to realize noise refrigeration of the reset resonator, where the QCR is driven through thermal noise [75, 24]. Channeling undesirable thermal noise into efficient cooling, this concept could reduce energy requirements of low-temperature quantum technological devices [76].

### 5.4 Two-tone Driving and Qubit Reset

Having experimentally examined the effect of the QCR on the reset resonator, we move on to studying its action on the qubit within the same circuit. Building on the results obtained from double-junction experiments in Sec. 4.2, we utilize the single-junction QCR sample in Fig. 5.6 (a), which was already measured in the previous section.

Due to the reduced coupling strength through the detuned reset resonator, we expect longer reset times for the dc-operation of the QCR compared to the direct-coupling experiment. To mitigate this, we apply a reset protocol combining elements of $f_{0-g_1}$ reset [8] and dc-biased QCR reset. This two-tone drive protocol [20, 21], theoretically described in Sec. 2.3.5, consists of a $f_{0-g_1}$ drive pulse that is assisted with an $e-f$ pump pulse to minimize the $e$-state population in the qubit. The $f_{0-g_1}$ drive frequency is chosen so that it creates an excitation in the reset resonator while the readout resonator remains unaffected. We improve upon the classic $f_{0-g_1}$
reset protocol by applying a dc bias to the QCR during the $f_0$-$g_1$ drive pulse to temporarily create strong dissipation in the reset resonator and absorb excitations caused by the $f_0$-$g_1$ pulse. This pulse sequence is shown in Fig. 5.6 (b).

As a first observation, the 6.6 µs decay time of the $g$-$e$ transition in the QCR-off state is longer than the 1.76 µs measured in Sec. 4.2 for direct capacitive coupling. This can partly be attributed to filtering of the reset resonator, but differences in fabrication technique must also be taken into account. With the reset pulse only consisting of a QCR dc bias and $V_{f_0g_1} = V_{e_f} = 0$, we see an accelerated decay of the $e$ and $f$-state in Fig. 5.7. This time, the trajectories follow an exponential behavior even at bias voltages near $eV = \Delta$.

In order to optimize the amplitudes of each drive tone, we calibrate the Rabi frequencies using time-domain measurements [8]. The drive amplitude corresponding to a specific Rabi frequency can be fitted using a Liouvillian model truncated at the third excited state. This model also provides an estimate of the reset resonator decay rates for different QCR amplitudes, yielding $1/\gamma^{\text{off}} = 221$ ns and $1/\gamma^{\text{on}} = 120$ ns. The $f_0$-$g_1$ drive frequency is calibrated experimentally to compensate for an amplitude-dependent Stark-shift. Four different sets of parameters $A, B, C, D$ are chosen and listed in Tab. 5.2. For reset purposes, we choose a conservative QCR amplitude of $V = 0.16$ mV to prevent any heating or excitation effects.

As the $f_0$-$g_1$ drive creates excitations outside of the two-level subspace, we also investigate the populations of higher qubit levels up to the third excited state using single-shot readout as described in Sec. 3.3. The measured distributions for different reset pulse lengths
Figure 5.7. Decay of the $e$ (a) and $f$-state (b) measured with averaged readout in the absence of a two-tone qubit drive. The dashed lines mark the QCR bias amplitudes $V=0$ mV, $V=0.16$ mV, and $V=0.22$ mV (gap voltage). (c) Cross-sections of the data with amplitudes corresponding to the dashed lines in (a) and (b), accompanied by fits of exponential decay with time constant $T_d$. Figure adapted from Publication VII.

Table 5.2. Parameter sets A–D for two-tone reset experiments. The off-state dynamics are represented by A, $f_0$–$g_1$ reset by B, QCR reset by C, and the combined QCR and $f_0$–$g_1$ protocol by D.

<table>
<thead>
<tr>
<th>Pulse amplitude (µV)</th>
<th>A</th>
<th>B</th>
<th>C</th>
<th>D</th>
</tr>
</thead>
<tbody>
<tr>
<td>$V_{QCR}$</td>
<td>0.0</td>
<td>0.0</td>
<td>160</td>
<td>160</td>
</tr>
<tr>
<td>$V_{f_0g_1}$</td>
<td>0.0</td>
<td>169.1</td>
<td>0.0</td>
<td>253.6</td>
</tr>
<tr>
<td>$V_{ef}$</td>
<td>0.0</td>
<td>10.6</td>
<td>0</td>
<td>16.9</td>
</tr>
</tbody>
</table>

are shown in Fig. 5.8 (a)–(f) and are fitted with a four-component GMM. This measurement follows the full pulse sequence in Fig. 5.6 (b) with drive amplitudes according to Tab. 5.2. Even though single-shot measurements allow us to extract individual populations of each qubit state, it is our goal to reduce all excitations, leading us to analyze the combined excited state population $P_{exc} = P_e + P_f + P_h$.

The dynamics of $P_{exc}$ are measured in Fig. 5.8 (g)–(l) for different drive amplitudes. Starting from thermal equilibrium, we observe a small rise in excited state population. For initial states $e$ and $f$, an accelerated decay is measured for parameter sets B, C, and D in comparison to A. The case C corresponds to an identical reset procedure to that studied in Sec. 4.2, but the decay is slower due to a smaller coupling strength through the reset resonator compared with a strong capacitive coupling. As expected, the fastest decay is produced by the parameter set D since it combines the increased decay rates induced by the QCR and by the two-tone drive. In this case, we reach a stabilization time around 500 ns, significantly faster than with a pure QCR bias pulse and 20 times faster than intrinsic excited state decay. Within this time range, the measured populations are in good agreement with master equation simulations.
Many reset protocols based on microwave drives only remove excitations from the first excited state [77, 9], but this experiment shows the simultaneous removal of higher excitations in the qubit. Whereas the two-tone drive by itself only acts on the $g$, $e$, and $f$ states, the QCR-induced dissipation affects all qubit states equally. The transmon temperature of 110 mK is relatively high, but largely unaffected by the QCR drive within the range chosen in Table 5.2, see the cases A and C in Fig 5.8 (g) and (h). An increased population of the excited state can be observed for the parameter sets B and D due to the two-tone drive, but this does not correspond to an increase in qubit temperature. Notably, the thermal equilibrium excitation, influenced by the cryogenic environment and quasiparticle generation in the QCR and coupled circuit [78], is lower than in similar experiments by Yoshioka et al. [21] and can possibly be reduced further by the implementation of quasiparticle barriers and Cu baths as described in Sec. 3.1. As a simplification, a dispersively coupled resonator with a QCR can be driven for reset and readout simultaneously, minimizing the circuit footprint and allowing higher scalability for multi-qubit systems and quantum processors.
5.5 Rapid Generation of Thermal States

As a last experiment, we want to achieve the exact opposite of the previous chapter: creating excitations in the transmon-resonator system with the QCR. In contrast to the QCR amplitude being limited by $\Delta/e$ in the two-tone experiment, we explore the bias range from $eV=\Delta$ to $5.5\Delta$ and analyze the population distribution up to the fifth excited state. We find these populations to be Boltzmann-distributed, corresponding to a Gibbs state of the qubit as shown in Fig 5.9 (d).

We use the measurement setup as described in Sec. 3.2 and Fig. 3.3, and once again employ single-shot readout. This allows us to monitor the qubit state distributions and employ a GMM to analyze the populations as shown in Fig. 5.9. Our fits of the populations to the closest Boltzmann distribution suggest that the $1\sigma$-demarcation of the $h$-state also includes measurements of the $i$ and $j$ state, so we treat it as the combined population of the third to fifth excited state.

Fig. 5.10 (a) shows the temperature corresponding to the fits of the Boltzmann distribution in Fig. 5.9 (d) for a fixed QCR pulse.
length of 100 ns. With increasing amplitude we see an increase in temperature, as is expected from Eq. (2.26). Within 100 ns we reach a temperature of 487 mK with a QCR pulse amplitude of 1.2 mV, the equivalent of $5.5\Delta/e$. The temperature fits become less accurate for higher QCR amplitudes, as higher temperatures populate higher qubit levels and the truncation error of the measurement becomes more significant.

We also compare time-domain measurements of different pulse amplitudes, finding an exponential behavior with time parameters in the range of 185 ns for $V=0.3$ mV, 80 ns for $V=0.6$ mV, and 109 ns for $V=1.2$ mV in Fig. 5.10 (b).

Several heating mechanism can be considered to contribute to the temperature increase in Fig. 5.10. The primary mechanism is photon- assisted tunneling causing the emission of photonic excitations into the circuit as depicted in Fig. 2.2 (c) [53]. Apart from photonic heating, tunneling processes in the QCR can generate quasiparticles [79, 80, 81], which can cause elevated temperatures even at zero bias, and their recombination can create ballistic phonons in the substrate, which can cause disturbances in the qubit [82, 83, 84].

Having experimentally realized the cooling of the resonator as well as the excitation-removal and thermal states in the qubit, we can combine these attributes to propose the use of the QCR as a two-way tunable environment in a potential quantum heat engine [26]. In this case, the QCR-resonator system could replace the hot and cold reservoir with a tunable environment, providing a large temperature tuning ratio. In combination with frequency-tunable qubit, this en-
ables the implementation of a quantum-Otto-cycle [85]. Alternatively, this system could be utilized for a minimal heat engine consisting of bandgap reservoirs coupled to a qubit, constituting a possible platform for the study of qubit-environment correlations [86, 87, 88].

It is important to note that the experiments in Chs. 5.3–5.5 have been somewhat limited by a QCR with a relatively high Dynes parameter in the range of $10^{-3}$. Both cooling of the resonator and excitation-removal in the qubit, as well as the temperature control of the transmon-resonator system can be further optimized by a smaller Dynes parameter, see Eq. (2.41). Dynes parameters in the range of $10^{-4}$ have been achieved experimentally in this context in Publication III and Ref. [21].
6. Summary and Conclusions

In this thesis, we explored the fascinating, yet challenging world of interactions between qubits, resonators, and the quantum-circuit refrigerator. Encompassing different realizations of these devices and circuits, comparisons, and combinations of different methods of qubit control were developed and put to the test in our lab, alongside theoretical and computational exercises to gain a deeper understanding of the behavior and dynamics of these quantum systems. At this point, we reflect on the advancements shown in this thesis as well as future considerations and the challenges ahead.

Our initial exploration of direct qubit reset with a double-junction QCR not only proved the feasibility of qubit reset with a QCR, but did so on an extremely competitive timescale below 10 ns. However, we found unexpected behavior for high pulse amplitudes, including even a temporary increase in excitation, leading us to choose more conservative parameters for our reset protocol. This highly promising experiment served as a proof-of-concept and inspired the study of more advanced qubit-resonator-QCR circuits as well as the conception of the single-junction QCR.

The single-junction QCR was first developed in a theoretical context in a QCR-resonator system, finding similar properties to its predecessor. Initially motivated by device simplification and more accurate bias control at the NIS junction, this study also sheds light on fascinating non-linear phenomena in high-impedance resonators, which were previously unexplored. Additionally, the single-junction QCR allowed us to investigate a chain of three coupled resonators, supporting the investigation of exceptional points and the associated system dynamics, the first study investigating a system with multiple QCRs and multiple circuit elements.

We then set our eyes on the highly versatile QCR-resonator-qubit circuit, aiming to experimentally realize a hybrid approach to qubit reset combining microwave-driven reset protocols with QCR-induced dissipation. Building on theoretical results, we developed the new
circuit using electromagnetic simulations and first delved into the resonator-QCR interaction and cooling properties through Fock-state measurements. We used different artificial heating methods to evaluate its cooling properties and discussed the implementation of noise-driven refrigeration.

Using single-shot readout we were able to resolve the distribution of qubit state populations including higher states outside of the two-level computational subspace and gained insight into their dynamics. This method was utilized for the experimental investigation of a hybrid reset protocol combining microwave-driven \( f_{0-g_1} \) reset with QCR-based dissipation in the resonator, removing excitations from the whole superconducting circuit. This method, albeit slower than direct reset, protects qubit coherence and allows for a longer qubit lifetime in the off-state.

Opening another avenue in the study of QCRs, the fast experimental generation of thermal states was described in the penultimate section. This ability, at first sight conflicting with our goal of qubit reset, is highly sought after in the field of quantum thermodynamics, for example for heat transport experiments in the quantum realm. We discussed the implementation of a possible quantum heat engine using the QCR as a two-way tunable bath.

As we set our eyes on future applications of this versatile device, the logical continuation of heating experiments is the QCR-driven quantum heat engine, which can potentially be realized experimentally with the same circuit as used in Publications I, VI, and VII. If successful, this would constitute the first quantum heat engine realized in a superconducting circuit [27], offering a unique opportunity in the fundamental research of quantum thermodynamics. On the trajectory of qubit reset, further optimization toward faster reset times, higher fidelities, and increased on/off ratios can be achieved through improved coupling strengths and, most importantly, the reduction of the Dynes parameter in the fabrication process. Within the QCR-resonator subsystem, the particular setting of a QCR coupled to a high-impedance resonator could potentially provide further applications and experiments based on its exhibition of non-linear phenomena.

This thesis provides a collection of theoretical and experimental adventures into the world of quantum-circuit refrigeration. Combining these results, we have found extreme versatility in this device and gained a deep understanding of its interaction with different circuits. In the conclusion of this thesis, we look ahead toward numerous applications of the QCR in circuit quantum electrodynamics and quantum thermodynamics far beyond qubit reset.


References


