

DOCTORAL THESIS RESEARCH PROPOSAL

Disentangling radioactive backgrounds in superconducting qubits with event-by-event detection

PHD PROGRAM IN ASTROPARTICLE PHYSICS: XXXVII CYCLE

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Abstract

✍ In the past twenty years, the greatest advances towards the realization of fault-tolerant quantum computers have been achieved by employing quantum bits, or qubits, based on superconducting circuits. However, the quantum states of superconducting qubits remain inherently fragile, as they are affected by several noise mechanisms due to unwanted couplings with the surrounding environment. Recently, increasing concerns for the impact of environmental radioactivity on quantum circuits have arisen: hitting particles, in fact, produce phonons in the chip substrate that can disturb multiple qubits simultaneously, complicating the implementation of error correction algorithms. This thesis deeply investigates how radiation sources affect superconducting qubits. The first step of this project was the realization of a Monte Carlo simulation to understand which radiation sources produce the highest rate of impacts in qubit chips and should, therefore, be taken into account for the implementation of mitigation strategies. Following that, two studies were performed, respectively on a fluxonium and on a transmon qubit. The first study aimed at understanding the impact of environmental radioactivity on the global decay time of the fluxonium qubit, concluding that other decay mechanisms dominate. The second study enabled the detection of individual particle interactions on a transmon chip, demonstrating that gamma radiation can induce qubit decays and establishing the transmon as a real-time radiation detector. Building on this result, a new firmware for the control and readout electronics was developed to enable multiplexed operation and the simultaneous measurement of up to four qubits. Having demonstrated real-time radiation detection with a single qubit, the multiplexed architecture allows the study of correlated errors across multiple qubits, a crucial step toward understanding radiation-induced error bursts in quantum processors. Initial tests of this firmware, performed by illuminating a chip with an LED, revealed correlated errors in three transmon qubits, paving the way for systematic investigations of radiation-induced correlated events. More broadly, by characterizing the transmon as a sensitive probe of substrate phonons, this work opens the door toward a new paradigm of phonon-mediated detection with superconducting qubits, a possibility that had previously been largely theoretical.

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Introduction

Over the past decades, the research on quantum technologies has made significant progress toward the realization of a fully-functioning quantum computer. Originally proposed by Yuri Manin [1], Paul Benioff [2], and Richard P. Feynman [3] to overcome the difficulties in simulating quantum systems on classical computers, quantum computers are characterized by intrinsic quantum behavior, i.e. they feature properties like superposition and entanglement specific to quantum systems.

Theoretical developments in quantum algorithms proved [4–7] that a quantum computer is capable of solving computational problems that cannot be solved efficiently by classical computers, and as of today the possible applications of quantum computing span a wide range of fields, including cryptography [8], machine learning [9], quantum simulations [10] and high-energy physics [11].

Despite the theoretical results, the realization of a quantum computer has proven to be much more difficult to achieve. The list of technologies that have been proposed over the years includes, among others, superconducting circuits [12], atomic ions confined by electric fields [13], semiconductors [14], photons [15] and nuclei [16], but for all of them there are complications concerning scalability, error rate and other requirements that have not been fully overcome yet. Nonetheless, the performance of quantum processors has improved dramatically in recent years, and in 2019 Google’s Sycamore processor was able to achieve in around 200 seconds a computational result that would have taken thousands of years for a classical supercomputer to obtain [17], confirming for the first time the theoretical expectations.

Since Google’s demonstration of quantum advantage in 2019, quantum processors have grown in complexity and reliability, but still do not satisfy all the requirements for most of their applications. In particular, their error rate remains too high to reliably execute complex quantum algorithms. These errors are mostly due to interactions with the surrounding environment, which cause the quantum state to evolve unpredictably. This phenomenon, known as *decoherence*, will be discussed in Chapter 2. As of today,

understanding decoherence mechanisms and mitigating them is one of the main research topics among scientists involved in quantum technologies.

The work of this thesis focuses on the impact of environmental radioactivity on the coherence of superconducting quantum bits. Although it is overall a minor contributor to qubit decoherence (as will be shown in Chapter 5), the multi-qubit errors it produces challenge error-correction algorithms and are, thus, a source of great concern.

The first part of this PhD project, reported in Chapter 4, was a systematic study of the different sources of radioactivity and their impact on superconducting qubit chips. Measurements of environmental radiation sources were performed in different experimental locations, while typical setup elements of superconducting qubit experiments were screened to estimate the activity of radioactive contaminants inside them. Building on these measurements, a Monte Carlo simulation was performed to assess, for each source, the rate of particle interactions in qubit chips.

Chapter 5, instead, describes a comparative study of the performance of a fluxonium qubit in different radiation conditions. In the experiment, the qubit was probed both when shielded against environmental radioactivity and when exposed to γ radiation sources, concluding that other decoherence mechanisms were playing a much greater role in determining the average coherence time of that qubit.

The study on the fluxonium qubit highlighted the need for a different measurement strategy for detecting particle interactions. This was developed in a second study, described in Chapter 6, where for the first time a transmon qubit was operated as a γ radiation detector, demonstrating not only that environmental γ -rays are a source of decoherence for superconducting qubits, but also that this technology could be used to develop a novel particle detector.

Finally, Chapter 7 describes the development and test of a new firmware for the detection of multi-qubit errors, an important step towards the development of chip designs more resistant to ionizing radiation. A first measurement of correlated errors induced by an LED was completed successfully, paving the way for future studies in that field.

Chapter 1

Superconducting Qubits

1.1 Quantum bits

The fundamental unit of information in a quantum computer is the quantum bit, or *qubit* [18]. A qubit is a two-level quantum system with $|0\rangle$ and $|1\rangle$ as its eigenstates, which serve as the quantum analogues of the classical bit states 0 and 1. The key difference between classical bits and qubits is that, whereas the classical bit can be set only to 0 or 1, the qubit can also be prepared in a *coherent superposition* of the states, i.e. a linear combination of $|0\rangle$ and $|1\rangle$:

$$|\psi\rangle = \alpha|0\rangle + \beta|1\rangle \quad (1.1)$$

Where α and β are complex numbers such that $|\alpha|^2 + |\beta|^2 = 1$. When measuring the qubit state $|\psi\rangle$, the outcome is $|0\rangle$ or $|1\rangle$ with probabilities $|\alpha|^2$ and $|\beta|^2$, respectively, but the role of the two coefficients goes beyond that: their relative phase enables interference effects and entanglement with other qubits, phenomena that have no classical analogue.

Since $|\alpha|^2 + |\beta|^2 = 1$, Eq.1.1 can be rewritten as [19]:

$$|\psi\rangle = e^{i\gamma} \left(\cos \frac{\vartheta}{2} |0\rangle + e^{i\varphi} \sin \frac{\vartheta}{2} |1\rangle \right) \quad (1.2)$$

Where γ , ϑ and φ are real numbers ranging from 0 to 2π . Quantum states are invariant under global phase transformations, so the factor $e^{i\gamma}$ can be neglected and Eq.1.2 becomes:

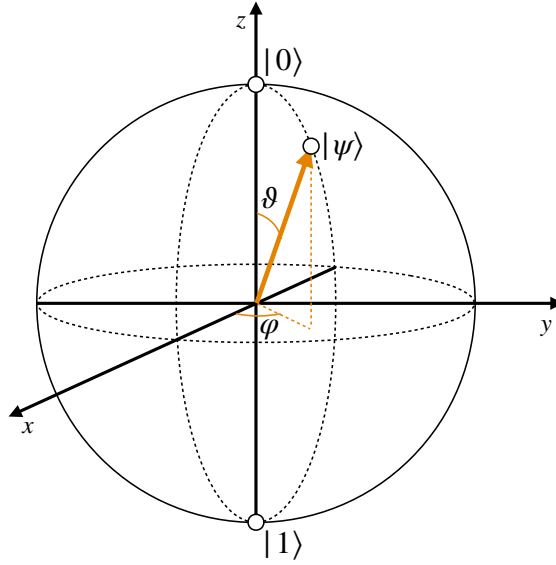


FIGURE 1.1: The Bloch sphere used to represent qubit states.

$$|\psi\rangle = \cos \frac{\vartheta}{2} |0\rangle + e^{i\varphi} \sin \frac{\vartheta}{2} |1\rangle \quad (1.3)$$

This representation allows mapping each qubit state into a point on the surface of a three-dimensional sphere with unit radius, known as the *Bloch sphere*. On this sphere the north pole corresponds to the eigenstate $|0\rangle$, while the south pole corresponds to the eigenstate $|1\rangle$. The angles ϑ and φ are then, respectively, the polar angle and the azimuthal angle. The z axis, being the one with the $|0\rangle$ and $|1\rangle$ states, is known as the *qubit quantization axis* or the *longitudinal axis*. The x and y axes, instead, are known as the *transverse axes*.

Due to the possibility of preparing superposition states, the manipulation of a qubit is much less trivial than that of a classical bit. In a classical computer, the basic single-bit operators, i.e., the identity and NOT *gates*, leave the state unchanged or flip it from 0 to 1 and vice versa. Other logic gates act on two or more bits, combining their values to perform more complex operations. In a qubit, instead, a *quantum logic gate* can change the qubit state to any superposition state. That means, any operator that manipulates the α and β coefficients of the superposition state $|\psi\rangle$ while preserving the normalization $|\alpha|^2 + |\beta|^2 = 1$ is a quantum logic gate. By representing the eigenstates $|0\rangle$ and $|1\rangle$ as orthogonal column vectors:

$$|0\rangle = \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad |1\rangle = \begin{pmatrix} 0 \\ 1 \end{pmatrix} \quad (1.4)$$

A generic superposition state can be written as:

$$|\psi\rangle = \begin{pmatrix} \alpha \\ \beta \end{pmatrix} \quad (1.5)$$

With norm $|\alpha|^2 + |\beta|^2 = 1$. As a consequence, any 2×2 operator U acting on $|\psi\rangle$ while preserving its norm is a single-qubit logic gate. Requiring the operator to preserve the norm of $|\psi\rangle$ is equivalent to asking the operator to be *unitary*, i.e. such that $U^\dagger U = U U^\dagger = I$.

1.2 Superconducting Qubits

One of the most promising ideas for the realization of reliable quantum computers consists in using superconducting circuits as qubits. Compared to other technologies, superconducting circuits have several advantages that favored their growth as one of the leading technologies for quantum computers: they offer very fast gate operations (down to nanoseconds) and rapidly improving gate fidelities; their fabrication and design allows to exploit technologies closely related to classical integrated microelectronics, including lithography, deposition, wiring, thus facilitating manufacturability and scalability; finally, their energy-level spectra are governed by the parameters of the circuit elements, making it possible to modify them by just changing the circuit design [20–22].

Over the last twenty-five years several realizations of a superconducting qubit have been proposed and tested, but almost all of them derive from two fundamental designs: the *charge qubit* and the *flux qubit* [23]. As their name suggest, in the first case the qubit state is determined by different charge configurations on a superconducting island, while in the latter the state is determined by the number of magnetic flux quanta threading a superconducting loop.

In both cases the dynamics is determined by the presence of one or more *Josephson junctions*, which are small dielectric junctions between two superconductors. For qubits deriving from charge quantization the Josephson junction is the link between the superconducting island and the rest of the circuit, with the tunneling of Cooper pairs through that determining the evolution of the qubit state. For qubits deriving from flux quantization, instead, the Josephson junction interrupts the superconducting loop and provide a phase drop that enables distinct quantized current states. To simplify, the junction's action is similar to the one of a valve, allowing flux quanta to enter or leave the loop.

The starting point for the mathematical description of superconducting qubits is the quantum LC resonator. LC resonators are among the simplest superconducting circuits

that can be realized, and they already exhibit, when some conditions are satisfied, quantum behavior, having a quantized energy spectrum. However, due to the difficulty in keeping them only in the first two energy states, they are unsuitable for working as qubits. This problem was solved by introducing the Josephson junction in the circuit, which provides a non-linear inductive element and breaks the harmonicity of the spectrum. The resulting device, known as the *Cooper-pair box*, was the first superconducting circuit successfully operated as a qubit. The Cooper-pair box will be the starting point for the derivation of the two qubit designs used in this thesis project, the *transmon* and the *fluxonium*. Finally, a brief overview of other superconducting qubit designs will be given.

1.2.1 The quantum LC resonator

An LC resonator (Fig.1.2a) is described by the Hamiltonian

$$H_{LC} = \frac{Q^2}{2C} + \frac{\Phi^2}{2L} \quad (1.6)$$

Where the first term is the energy term of the capacitor, determined by the charge Q and the capacitance C , and the second term is the energy term of the inductor, determined by the magnetic flux Φ through the inductor and the inductance L .

The inductive term can be rewritten by introducing the resonant frequency $\omega_r = 1/\sqrt{LC}$, obtaining:

$$H_{LC} = \frac{Q^2}{2C} + \frac{1}{2}C\omega_r^2\Phi^2 \quad (1.7)$$

In this form, the Hamiltonian is equivalent to the one of the harmonic oscillator:

$$H_{osc} = \frac{p^2}{2m} + \frac{1}{2}m\omega^2x^2 \quad (1.8)$$

Giving a hint of how quantization can be achieved. In the quantum harmonic oscillator the momentum p and the position x are promoted to operators with commutation relation $[\hat{x}, \hat{p}] = i\hbar$. Following the same path, the LC resonator can be quantized by promoting Q and Φ to operators and assigning them the commutation relation:

$$[\hat{\Phi}, \hat{Q}] = i\hbar \quad (1.9)$$

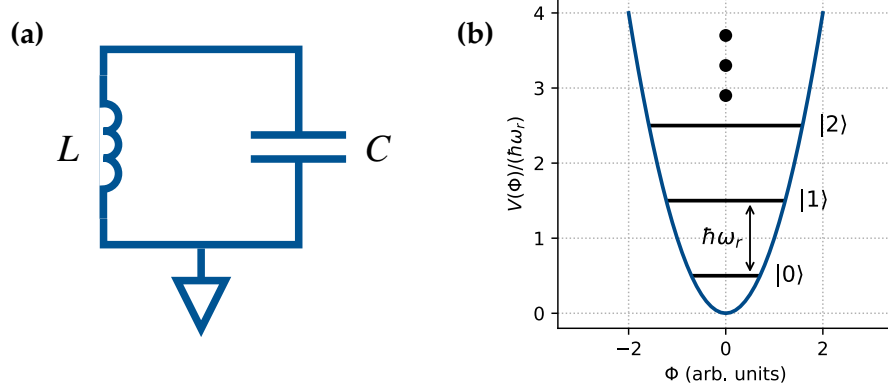


FIGURE 1.2: (a) Circuit diagram of a LC resonator and (b) its energy spectrum.

The analogy is not over, though. As in the quantum harmonic oscillator, also $\hat{\Phi}$ and $\hat{Q}$ can be rewritten in terms of *annihilation* and *creation* operators, namely $\hat{a}$ and $\hat{a}^\dagger$. The two sets of operators are related by the equations:

$$\hat{\Phi} = \sqrt{\frac{\hbar}{2\omega_r C}} (\hat{a}^\dagger + \hat{a}), \quad \hat{Q} = i\sqrt{\frac{\hbar\omega_r C}{2}} (\hat{a}^\dagger - \hat{a}) \quad (1.10)$$

Or, by introducing the *characteristic impedance* $Z_r = \sqrt{L/C}$:

$$\hat{\Phi} = \sqrt{\frac{\hbar Z_r}{2}} (\hat{a}^\dagger + \hat{a}), \quad \hat{Q} = i\sqrt{\frac{\hbar}{2Z_r}} (\hat{a}^\dagger - \hat{a}) \quad (1.11)$$

The operator $\hat{a}^\dagger$ creates a photon with energy $\hbar\omega_r$, while the operator $\hat{a}$ destroys one. By using $\hat{a}$ and $\hat{a}^\dagger$, instead of $\hat{\Phi}$ and $\hat{Q}$, the Hamiltonian becomes:

$$\hat{H}_{LC} = \hbar\omega_r (\hat{a}^\dagger \hat{a} + 1/2) \quad (1.12)$$

The eigenstates of the Hamiltonian satisfy the relation $\hat{a}^\dagger \hat{a} |n\rangle = n |n\rangle$, and because of that the $\hat{a}^\dagger \hat{a}$ is also known as *number* operator. Fig.1.2b shows the harmonic potential of the quantum LC resonator with its energy spectrum.

Having obtained a quantum description of the LC resonator, it is straightforward to wonder whether this quantum behavior can actually be observed. In order to do that, at least two requirements have to be satisfied:

1. The broadening of the spectral lines must be significantly smaller than the separation between levels, otherwise the overlaps would make it impossible to resolve the energy levels;

2. The energy separation between $|n\rangle$ and $|n+1\rangle$ must be larger than the thermal energy $k_B T$, otherwise the dynamics of the system would be dominated by spontaneous transitions between energy levels.

Both of these conditions are actually satisfied by present-day resonators. The broadening of the spectral lines is determined by the *quality factor*, which is defined as the ratio between the resonant frequency ω_r and the full width at half maximum (FWHM) κ of the resonance:

$$Q = \frac{\omega_r}{\kappa} \quad (1.13)$$

So the higher the quality factor, the lower the broadening. Currently, superconducting resonators have quality factors that can go beyond 10^6 [24], meaning that even if the separation between energy levels is not dramatically high it is possible to correctly resolve them.

In typical setups, instead, the energy separation is high, $\omega_r = E/\hbar = O(\text{GHz})$ and higher, because of the second condition: with such separation it is possible to observe a very clear spectrum, and thermally populate only the lowest energy states, by cooling down the circuit to $T \sim 10 \text{ mK}$, a temperature that can be reached by using a $^3\text{He} - ^4\text{He}$ dilution refrigerator [25].

The reason why LC resonators cannot be operated as qubits is the harmonic spectrum itself: since the energy separation is the same between $|0\rangle$ and $|1\rangle$, $|1\rangle$ and $|2\rangle$, and so on for all the energy states, it is practically impossible to keep the system in a subset of its spectrum with only two levels. Every time photons are sent to the system, trying to excite it from $|0\rangle$ to $|1\rangle$, it is likely to excite it also to higher energy states, an error that is known as *state leakage*.

The solution to this problem consisted in modifying the circuit making its spectrum anharmonic: in this way, the energy required to excite the system from $|0\rangle$ to $|1\rangle$ would be different from the energy required for exciting it from $|2\rangle$ to $|2\rangle$. Considering microwave pulses, which are the most common way to operate such superconducting circuits, different energies mean different photon frequencies, damping the possibility to have state leakage. As will be described in the next section, the anharmonicity will be achieved by replacing the inductor with a Josephson junction.

1.2.2 From the Cooper-pair box to the transmon qubit

In 1962, B. Josephson theoretically demonstrated [26] that a dissipationless supercurrent can flow between two superconductors separated by a thin insulating barrier - typically an oxide layer only a few nanometers thick. The magnitude of this supercurrent is given by

$$I = I_c \sin \varphi \quad (1.14)$$

Where φ is the phase difference between the wavefunctions describing the superconducting condensates on the two sides of the junction, while I_c is known as the *critical current*, i.e. the maximum current achievable before the junction switches to a resistive regime. Josephson also predicted that, if an external potential V is applied across the junction, the time evolution of φ will depend on that potential through the equation:

$$\frac{d\varphi}{dt} = \frac{2e}{\hbar} V \quad (1.15)$$

To connect this expression with a circuit description, it is convenient to introduce the concept of *branch flux*, defined as the time integral of the voltage across the junction:

$$\Phi(t) = \int_{-\infty}^t V(t') dt' \quad (1.16)$$

The branch flux generalizes the notion of magnetic flux; in general it is a circuit variable, and only in the case of a closed loop the sum of branch fluxes correspond to the physical magnetic flux threading the loop [27]. With that the phase φ becomes:

$$\varphi(t) = \frac{2e}{\hbar} \int V(t') dt' = \frac{2e}{\hbar} \Phi(t) = 2\pi \frac{\Phi(t)}{\Phi_0} \quad (1.17)$$

Where $\Phi_0 = h/2e$ is the flux quantum. Eq.1.14 can then be rewritten as:

$$I = I_c \sin \left(2\pi \frac{\Phi}{\Phi_0} \right) \quad (1.18)$$

From which the *Josephson inductance* can be derived:

$$L_J(\Phi) = \left(\frac{\partial I}{\partial \Phi} \right)^{-1} = \frac{\Phi_0}{2\pi I_c} \frac{1}{\cos(2\pi \Phi / \Phi_0)} \quad (1.19)$$

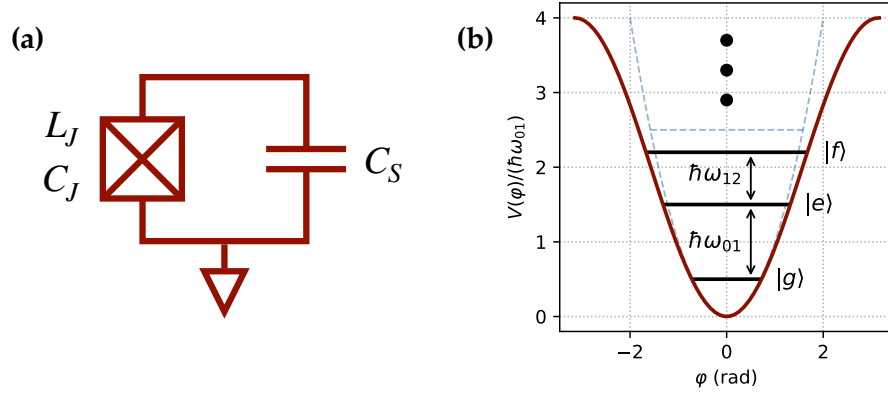


FIGURE 1.3: (a) Circuit diagram of a superconducting qubit and (b) its energy spectrum, compared with the one of the harmonic oscillator.

The Josephson inductance differs from the geometric inductance of the LC resonator described in the previous section by varying with the magnitude of the flux. As a consequence, the behavior of a Josephson junction is akin to the one of a non-linear inductor. Following that, also the energy stored in the junction depends on the flux, and is given by:

$$E = \int V(t)I(t)dt = \int \left(\frac{d\Phi}{dt} \right) I_c \sin \left(\frac{2\pi}{\Phi_0} \Phi \right) = -E_J \cos \left(\frac{2\pi}{\Phi_0} \Phi \right) \quad (1.20)$$

The term $E_J = I_c \Phi_0 / 2\pi$ is known as the *Josephson energy*.

By replacing the inductor in Fig.1.2a with a Josephson junction, as in Fig.1.3a, the circuit changes significantly. It now features a superconducting island delimited by the Josephson junction and the shunting capacitance. Also, its Hamiltonian differs from the one of the harmonic oscillator (see Eq.1.7), becoming, after promoting again Q and Φ to operators:

$$\hat{H} = \frac{\hat{Q}^2}{2C} - E_J \cos \left(\frac{2\pi}{\Phi_0} \hat{\Phi} \right) \quad (1.21)$$

Where $\hat{Q}$ is the excess charge in the island. That charge can be expressed in terms of Cooper pairs by replacing $\hat{Q}$ with the charge number operator $\hat{n} = \hat{Q}/2e$. By substituting also $\hat{\varphi} = (2\pi/\Phi_0)\hat{\Phi}$, such that $[\hat{\varphi}, \hat{n}] = i$, and introducing the charging energy $E_C = e^2/2C$, the Hamiltonian becomes:

$$\hat{H} = 4E_C \hat{n}^2 - E_J \cos \hat{\varphi} \quad (1.22)$$

The capacitance considered in this equation, though, is not only the shunting capacitance, but also the intrinsic capacitance of the Josephson junction: since it is made by two superconductors facing each other, separated by an insulating barrier, it also acts as a small capacitor.

Often, the superconducting island is coupled to a control electrode through a capacitance C_g , enabling the application of an external gate voltage V_g . The presence of the gate capacitance is crucial: by tuning V_g , it is possible to manipulate the charge configuration on the island and affect the qubit behavior. Mathematically, the gate voltage produces a gate charge:

$$n_g = \frac{C_g V_g}{2e} \quad (1.23)$$

That can be included in the Hamiltonian as:

$$\hat{H} = 4E_C(\hat{n} - n_g)^2 - E_J \cos \hat{\varphi} \quad (1.24)$$

This Hamiltonian can be diagonalized numerically to derive the energy spectrum. Due to the non-linearity of the Josephson junction, this spectrum is not harmonic anymore, having unequal energy separation between the levels (see Fig.1.3b).

The dynamics of this system, as well as its *anharmonicity*, depend on the E_J/E_C ratio. Earliest designs [28, 29], known as *Cooper-pair boxes*, worked in the $E_J < E_C$ regime, where the $4E_C(\hat{n} - n_g)^2$ term of the Hamiltonian is dominant. In this regime the eigenstates of the system closely resemble charge-number states $|n\rangle$. This is the reason why these qubits are known also as *charge qubits*.

The main flaw of these designs is their sensitivity to fluctuations of the charge offset n_g , what is known as *charge noise*. Since the value of n_g affects the energy separation between $|0\rangle$ and $|1\rangle$, and thus the frequency ω_{01} of the photons sent to manipulate the qubit state, charge noise ends up being a source of errors that cannot be neglected.

To overcome this limitation, the most common solution is to work in the *transmon* regime, where $E_J \gg E_C$ [30]. Here the Josephson term of the Hamiltonian dominates over the charge one, resulting in a lower sensitivity to charge noise due to the lower contribution of n_g . This is obtained by shunting the Josephson junction with a large capacitance, such that C increases and, as a consequence, E_C drops. However, the charge term cannot be reduced indefinitely, since doing that comes at the price of reducing also the anharmonicity of the system, making it harder to selectively address the computational transition $|0\rangle \leftrightarrow |1\rangle$ without populating higher excited states. For this

reason, transmon qubits are typically designed with E_J/E_C ratio ranging from 20 to 80, which is a reasonable compromise between protecting the qubit against charge noise and preserving the anharmonicity.

The relation between anharmonicity and the E_J/E_C ratio can be derived by Taylor expanding the cosine term around 0. This is justified by the fact that transmon qubits are operated much below the critical current I_c , meaning that (see Eq. 1.14) also the magnitude of φ has to be small. From the Taylor expansion the Josephson term of the Hamiltonian becomes:

$$E_J \cos(\hat{\varphi}) = -\frac{1}{2}E_J\hat{\varphi}^2 + \frac{1}{24}E_J\hat{\varphi}^4 + O(\hat{\varphi}^6) \quad (1.25)$$

And the Hamiltonian in equation 1.24 can be rewritten as:

$$\hat{H} \approx 4E_C\hat{n}^2 + \frac{1}{2}E_J\hat{\varphi}^2 - \frac{1}{24}E_J\hat{\varphi}^4 \quad (1.26)$$

Where n_g has been dropped being negligible. The first two terms of the Hamiltonian are analogous to the kinetic and potential terms of the harmonic oscillator, while the third term breaks the harmonic structure of the energy levels. In analogy with the quantum harmonic oscillator, it is possible to introduce two ladder operators, $\hat{b}$ and $\hat{b}^\dagger$, which are related to $\hat{n}$ and $\hat{\varphi}$ by:

$$\hat{\varphi} = \left(\frac{2E_C}{E_J}\right)^{1/4} (\hat{b}^\dagger + \hat{b}), \quad \hat{n} = \frac{i}{2} \left(\frac{E_J}{2E_C}\right)^{1/4} (\hat{b}^\dagger - \hat{b}) \quad (1.27)$$

In terms of these operators, the Hamiltonian becomes:

$$\hat{H} = \sqrt{8E_CE_J} \left(\hat{b}^\dagger \hat{b} + \frac{1}{2} \right) - \frac{E_C}{12} (\hat{b}^\dagger + \hat{b})^4 \quad (1.28)$$

The contribution of the $(\hat{b}^\dagger + \hat{b})^4$ term can be estimated perturbatively. At first order, it will shift the energy of the n -th state by:

$$\Delta E_n = -\frac{E_C}{12} \langle n | (\hat{b}^\dagger + \hat{b})^4 | n \rangle = -\frac{E_C}{12} (6n^2 + 6n + 3) \quad (1.29)$$

The energy separation between $|n\rangle$ and $|n+1\rangle$ is therefore given by:

$$E_{n+1} - E_n = \sqrt{8E_CE_J} + \Delta E_{n+1} - \Delta E_n = \sqrt{8E_CE_J} - E_C(n+1) \quad (1.30)$$

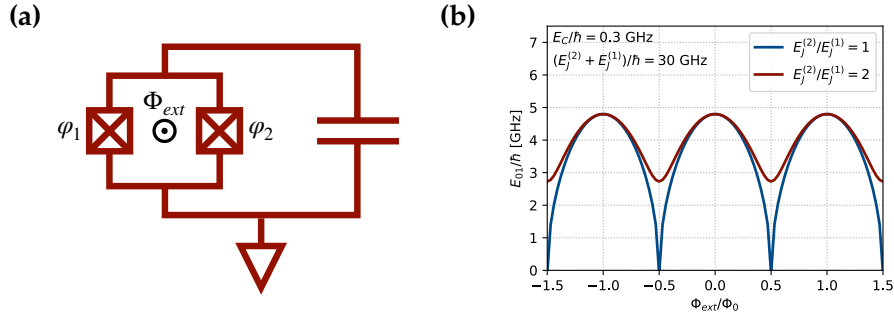


FIGURE 1.4: (a) Circuit diagram of a flux-tunable transmon qubit and (b) its transition energy as a function of the external flux Φ_{ext} for different $E_J^{(2)}/E_J^{(1)}$ ratios.

Restricting to the $|0\rangle$ and $|1\rangle$ levels, the anharmonicity of the system is then given by $\alpha = -E_C$. As anticipated, the lower is the value of E_C , the lower is the anharmonicity of the qubit. The design of a transmon qubit therefore involves finding a balance between decreasing E_C , which protects the qubit against charge noise, and preserving a value of α large enough to avoid state leakage when manipulating the qubit.

From the last equation it is also possible to derive the transition energy between $|0\rangle$ and $|1\rangle$, obtaining:

$$E_{01} = \sqrt{8E_CE_J} - E_C \quad (1.31)$$

Typical transmons have $E_{01}/\hbar = O(\text{GHz})$, large enough to prevent spontaneous excitations when cooled down to 10 mK. Typical anharmonicity values are instead $|\alpha/\hbar| \sim 100 - 400$ MHz.

1.2.3 The flux-tunable transmon qubit

The transmon qubit considered until now has a transition energy that is determined solely by the circuit parameters E_C and E_J . Its E_{01} is therefore fixed once the qubit has been fabricated. Many complex algorithms, though, require this value to be tunable over the computation [31–34]. An alternative design, that takes into account this necessity, consists in replacing the Josephson junction with two parallel ones, making what is known as a *Superconducting Quantum Interference Device* or SQUID [35] (see Fig.1.4a).

With this design, the qubit now has a superconducting loop interrupted by the two junctions, making it sensitive to the external magnetic flux Φ_{ext} threading that loop. This is due to flux quantization, that imposes the condition [36]:

$$\hat{\varphi}_1 - \hat{\varphi}_2 = 2\pi \frac{\Phi_{ext}}{\Phi_0} \quad (1.32)$$

Where $\hat{\varphi}_1$ and $\hat{\varphi}_2$ are the superconducting phase differences of the two junctions. The Hamiltonian of this new system, which is:

$$\hat{H} = 4E_C \hat{n}^2 - E_J^{(1)} \cos \hat{\varphi}_1 - E_J^{(2)} \cos \hat{\varphi}_2 \quad (1.33)$$

Can then be rewritten by replacing $\hat{\varphi}_1$ and $\hat{\varphi}_2$ with the average phase difference:

$$\hat{\varphi} = \frac{\hat{\varphi}_1 + \hat{\varphi}_2}{2} \quad (1.34)$$

The result is a shifted cosine term:

$$\hat{H} = 4E_C \hat{n}^2 - E_J(\Phi_{ext}) \cos(\hat{\varphi} - \varphi_0) \quad (1.35)$$

Where the Josephson energy E_J now depends on the external magnetic flux Φ_{ext} and is given by:

$$E_J(\Phi_{ext}) = [E_J^{(2)} + E_J^{(1)}] \cos\left(\frac{\pi\Phi_{ext}}{\Phi_0}\right) \sqrt{1 + d^2 \tan^2\left(\frac{\pi\Phi_{ext}}{\Phi_0}\right)} \quad (1.36)$$

With d being the junction asymmetry:

$$d = \frac{E_J^{(2)} - E_J^{(1)}}{E_J^{(2)} + E_J^{(1)}} \quad (1.37)$$

As a consequence, varying the external magnetic flux Φ_{ext} will change the transition energy E_{01} , that now is given by:

$$E_{01}(\Phi_{ext}) = \sqrt{8E_C |E_J(\Phi_{ext})|} - E_C \quad (1.38)$$

And will oscillate between two values (see Fig.1.4b). When $E_J^{(1)} = E_J^{(2)}$, so the junction asymmetry d is 0, $E_J \propto \cos(\pi\Phi_{ext}/\Phi_0)$ and E_{01} oscillates between 0 and $\sqrt{16E_CE_J^{(1)}} - E_C$. As $|d|$ increases, one of the two junctions becomes dominant, reducing the amplitude of the E_{01} oscillations. In the extreme case, with $|d| = 1$, only one junction affects the qubit, and the flux tunability vanishes.

As protecting qubits from charge noise came at the price of sacrificing anharmonicity, also for flux-tunability there is a trade-off to consider. In fact, with E_{01} being sensitive

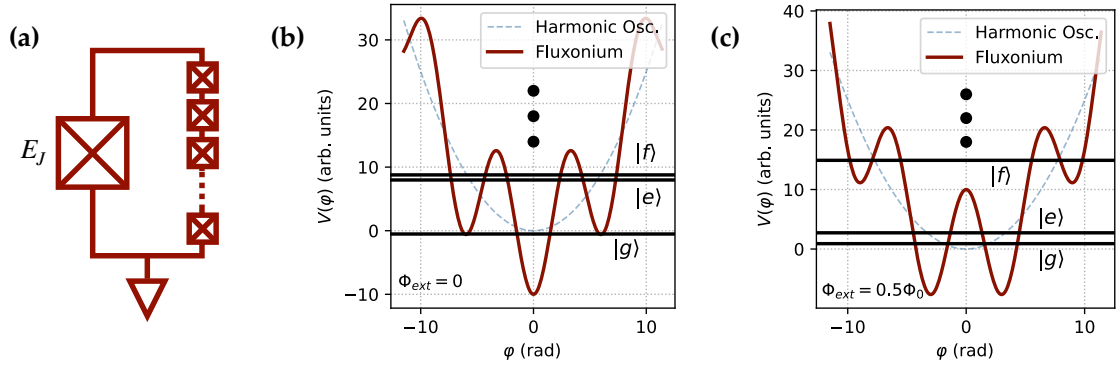


FIGURE 1.5: (a) Circuit diagram of a fluxonium qubit, (b) its energy spectrum for $\Phi_{ext} = 0$ and (c) for $\Phi_{ext} = 0.5\Phi_0$.

to the magnetic flux threading the loop, fluctuations of this flux represent a new source of noise to be taken into account. This is known as *flux noise*, and plays a role also in the next qubit design relevant for the work of this thesis.

1.2.4 The fluxonium qubit

As discussed in section 1.2.2, the small anharmonicity remains the main drawback of transmon qubits. Because of that, despite the transmon is still the most common design used for prototypes of quantum computers, other designs have been studied, with the aim of abating the charge noise affecting Cooper-pair boxes without undermining anharmonicity. One of the most promising ideas consists in dropping the superconducting island altogether, replacing it with a superconducting loop interrupted by one or more Josephson junctions. In this design, the quantum state is determined by the number of magnetic flux quanta trapped inside the loop. Following this idea, Manucharyan *et al.* introduced the *fluxonium* qubit, where the Josephson junction is shunted by a large inductance [37]. The high inductance required by this qubit was achieved by realizing a long array of Josephson junctions (see Fig.1.5a), typically tens or hundreds in modern prototypes.

Similar to what has been done for the flux-tunable transmon, the Hamiltonian of this device can be written as:

$$\hat{H} = 4E_C\hat{n} - E_J \cos(\hat{\varphi}) - \sum_{n=1}^N \gamma_n E_J \cos(\hat{\varphi}_n) \quad (1.39)$$

With $\gamma_n = E_J^{(n)}/E_J$. Also here, flux quantization has to be imposed, obtaining:

$$\hat{\varphi} - \sum_{n=1}^N \hat{\varphi}_n = 2\pi \frac{\Phi_{ext}}{\Phi_0} \quad (1.40)$$

For simplicity, it is possible to rename $\sum_{n=1}^N \hat{\varphi}_n = \hat{\phi}$ and $2\pi\Phi_{ext}/\Phi_0 = \varphi_{ext}$. Under the hypothesis that the Josephson junctions of the array are all equal, the Hamiltonian becomes:

$$\hat{H} = 4E_C \hat{n} - E_J \cos(\hat{\phi} + \varphi_{ext}) - N\gamma E_J \cos\left(\frac{\hat{\phi}}{N}\right) \quad (1.41)$$

Being N large, the last cosine can be Taylor expanded:

$$\hat{H} = 4E_C \hat{n} - E_J \cos(\hat{\phi} + \varphi_{ext}) + \frac{1}{2} \frac{\gamma E_J}{N} \hat{\phi}^2 \quad (1.42)$$

The last term of the Hamiltonian is an inductive one, so by renaming $(\gamma/N)E_J = E_L$ the fluxonium Hamiltonian becomes:

$$\hat{H} = 4E_C \hat{n} - E_J \cos(\hat{\phi} + \varphi_{ext}) + \frac{1}{2} E_L \hat{\phi}^2 \quad (1.43)$$

The potential well, as seen in Fig.1.5, is a parabola modulated with a cosine term. Depending on the ratios between E_C , E_J and E_L two different designs are possible for fluxonium qubits, which will affect their potential and, as a consequence, their energy spectrum [38]:

- *light fluxonium*, where $E_L \ll E_J < E_C$, and the
- *heavy fluxonium*, where $E_C \approx E_L \ll E_J$.

For both regimes, the cosine term is too big to be treated as a perturbation, and the fluxonium energy levels are usually computed numerically (Fig.1.5b-c). In both cases, it is possible to achieve very high anharmonicity, up to $O(\text{GHz})$, overcoming the main limitation of transmon qubits.

The choice of relying on light or on heavy fluxonium is guided by the type of noise from which the qubit should be protected most. In a heavy fluxonium, the high E_J/E_C ratio strongly inhibits transitions from $|e\rangle$ to $|g\rangle$ and vice versa, allowing for long lifetimes of the qubit eigenstates (this is known as T_1 and will be described in section 2.3). At the same time, though, ω_{01} is very sensitive to the external flux Φ_{ext} (Fig.1.6, red line),

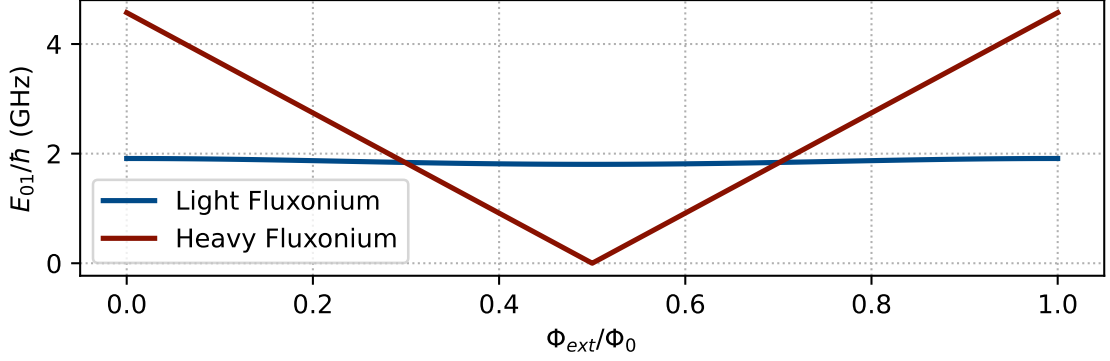


FIGURE 1.6: Transition energy, as a function of the external magnetic flux Φ_{ext} , for light and heavy fluxonium qubits. Values of E_C , E_J and E_L are taken from [39].

varying from few MHz up to GHz with different flux biases, making heavy fluxoniums highly sensitive to flux noise [39, 40].

In summary, experiments on fluxonium qubits showed their potential in overcoming the limitations of transmon qubits, especially regarding anharmonicity. However, their development is still at the beginning, and at this stage it is still impossible to protect fluxoniums both against unwanted state transitions and flux noise. As of today, research on both transmon and fluxonium designs is still ongoing, and it is too soon to claim which one will grant the best results.

1.2.5 Other designs for superconducting qubits

As mentioned earlier, transmon and fluxonium are only two of the many designs for superconducting qubits that have been proposed over the years. Other designs were not discussed in this chapter because all the devices used for this thesis project featured transmon or fluxonium qubits. This section provides a brief review of these other designs.

Prior to the transmon qubit, the *quantronium* qubit was proposed to improve the performance of Cooper-pair boxes [41]. Its architecture still relied on a small superconducting island, which was defined by two nominally identical Josephson junctions in series, forming a split Cooper-pair box embedded in a superconducting loop. A larger junction was then inserted in the loop to provide readout capabilities through the measurement of its switching current. The loop itself was then shunted with an external capacitance. A fundamental innovation introduced with quantronium was the possibility to operate it in a *sweet spot*, i.e. a regime such that the qubit was first-order insensitive both to charge and flux noise.

In the quest for abating charge noise, the transmon qubit was not the first qubit proposed relying on the superconducting phase $\hat{\varphi}$ as the relevant quantum variable. Before that,

a superconducting qubit based on phase quantization, and called because of that *phase* qubit, has also been proposed [42]. In this qubit the E_J/E_C ratio is significantly higher than in transmons, with typical values greater than 10^4 . In transmons, a ratio like that would make the anharmonicity negligible, but here the anharmonicity is preserved thanks to a DC current bias that distorts the potential well [43]. Distorting the potential well also changes the energy separation between levels, making it possible to tune these qubits by changing the DC bias, in a similar way to what is done in flux-tunable transmons with external magnetic fields.

Concerning qubits based on flux quantization, the first design to be proposed was the eponymous *flux* qubit, where the superconducting loop is interrupted by typically three Josephson junctions, engineered to create a double-well potential whose minima correspond to persistent currents circulating in opposite directions [44, 45]. At a flux bias of $\Phi_0/2$, coherent tunneling between these states forms the qubit levels. Later improvements introduced large shunt capacitors to reduce sensitivity to charge noise and enhance coherence [46, 47].

Sharing the same design principle as the fluxonium qubit are also the *blochonium* qubit [48] and the *bifluxon* qubit [49]. Blochonium qubits have circuit parameters similar to those of a light fluxonium qubits, with the E_J/E_C ratio small enough to make the charge term dominant. In this configuration Eq. 1.43 resembles the one of a particle in a weak periodic potential, like an electron in a crystal. Following this analogy, just as in crystals there are quasimomentum Bloch bands, with electrons undergoing transitions between them, here *quasicharge* Bloch bands can be defined, each associated with a different qubit state.

The design of Bifluxon qubits aims at protecting them both from charge and flux noise by making the parity of fluxons in the superconducting loop, i.e. whether the number of fluxons is even or odd, a conserved quantity. To do that, the loop has a superinductor shunting two identical Josephson junctions in series. The two junctions form a Cooper-pair box, which offset charge n_g can be tuned with an external potential. When $n_g \approx 0.5$, the parity of fluxons in the loop is preserved, so $|g\rangle \leftrightarrow |e\rangle$ transitions, which would change the number of fluxons from even to odd and vice versa, are suppressed.

Finally, another possible evolution of the fluxonium qubit is the $0-\pi$ qubit [50, 51], where the circuit is designed in such a way that the wavefunctions of $|0\rangle$ and $|1\rangle$ are not overlapping, so to protect the qubit from spontaneous $|0\rangle \leftrightarrow |1\rangle$ transitions.

Chapter 2

Control, readout and decoherence of superconducting qubits

2.1 Qubit control

2.1.1 Qubit interaction with an external EM field

A fundamental aspect of quantum computation is the coherent manipulation of the qubit state. When working with superconducting qubits, this process, commonly known as qubit driving, is achieved by using microwave pulses. These pulses are applied to the qubit through a coupling capacitance C_g , as shown in Fig.2.1a.

As discussed earlier for the Cooper-pair box (Section 1.2.2), an electrostatic potential applied to the superconducting island produces an offset charge in the qubit Hamiltonian. This also holds in this case, and it can be proven [21] that the Hamiltonian of this system, before quantization, is:

$$H = \frac{[Q - C_g V_g(t)]^2}{2(C_\Sigma + C_g)} - E_J \cos \varphi \quad (2.1)$$

Where C_g is the coupling capacitance and C_Σ is the sum of the shunting capacitance C_S and the intrinsic capacitance of the Josephson junction C_J . This can be rewritten as:

$$H = \underbrace{\frac{Q^2}{2(C_\Sigma + C_g)}}_{H_0} - E_J \cos \varphi - \underbrace{\frac{C_g}{C_\Sigma + C_g} Q V_g(t)}_{H_d} \quad (2.2)$$

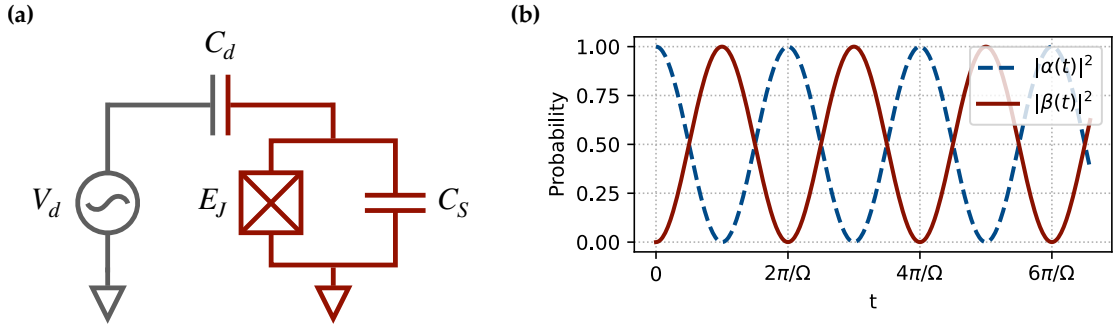


FIGURE 2.1: (a) Circuit diagram of a qubit coupled to a source V_d and (b) the Rabi oscillations of state probabilities $|\alpha(t)|^2$ and $|\beta(t)|^2$.

Where the term $(C_g^2 V_g^2)/[2(C_\Sigma + C_g)]$ has been dropped since it is just an offset that does not affect the dynamics of the system. The Hamiltonian just derived is the sum of the qubit Hamiltonian already encountered in Section 1.2.2, that will be named H_0 , and a perturbation term H_d .

Similar to what has been done in [22], Q and φ can be promoted to operators and substituted with the ladder operators $\hat{b}$ and $\hat{b}^\dagger$ (see Eq.1.27). The perturbation term then becomes:

$$\hat{H}_d = -i \frac{C_g}{C_\Sigma + C_g} V_g(t) Q_0 (\hat{b} - \hat{b}^\dagger) \quad (2.3)$$

Where all the constant terms in front of $(\hat{b} - \hat{b}^\dagger)$ have been grouped in Q_0 . Computationally, qubits are operated only in the $|0\rangle$ and $|1\rangle$ states, so they can be approximated as two-level systems. In this approximation, $\hat{b}$ and $\hat{b}^\dagger$ can be rewritten in terms of the Pauli matrices:

$$\hat{b} \rightarrow \hat{\sigma}_- = \frac{1}{2}(\hat{\sigma}_x - i\hat{\sigma}_y) \quad \hat{b}^\dagger \rightarrow \hat{\sigma}_+ = \frac{1}{2}(\hat{\sigma}_x + i\hat{\sigma}_y) \quad (2.4)$$

Obtaining:

$$\hat{H}_d = -i \frac{C_g}{C_\Sigma + C_g} V_g(t) Q_0 (\hat{\sigma}_- - \hat{\sigma}_+) \quad (2.5)$$

If the generator is producing a signal with frequency ω and amplitude V_g , the drive pulse can be written as:

$$V_g(t) = V_g \cos(\omega t) = \frac{V_g}{2} (e^{i\omega t} + e^{-i\omega t}) \quad (2.6)$$

Resulting in:

$$\hat{H}_d = -i \frac{C_g}{C_\Sigma + C_g} Q_0 \frac{V_g}{2} (e^{i\omega t} + e^{-i\omega t}) (\hat{\sigma}_- - \hat{\sigma}_+) \quad (2.7)$$

Before moving forward, the time evolution of the $\hat{\sigma}_-$ and $\hat{\sigma}_+$ operators also has to be considered. In the interaction picture, this is determined by the unperturbed Hamiltonian $\hat{H}_0$:

$$\frac{d\hat{\sigma}_-}{dt} = \frac{1}{i\hbar} [\hat{\sigma}_-, \hat{H}_0] \quad \frac{d\hat{\sigma}_+}{dt} = \frac{1}{i\hbar} [\hat{\sigma}_+, \hat{H}_0] \quad (2.8)$$

Using the same two-level approximation used for $\hat{H}_d$, the unperturbed Hamiltonian can be rewritten as:

$$\hat{H}_0 = \frac{\hbar\omega_{01}}{2} \hat{\sigma}_z \quad (2.9)$$

Where $\omega_{01} = (E_1 - E_0)/\hbar$ is the qubit frequency. Substituting eq. 2.9 into 2.8 and using the commutation relation $[\hat{\sigma}_j, \hat{\sigma}_k] = 2i\varepsilon_{jkl}\hat{\sigma}_l$, where ε_{jkl} is the Levi-Civita symbol, two relations are obtained:

$$\frac{d\hat{\sigma}_-}{dt} = -i\omega_{01}\hat{\sigma}_- \quad \frac{d\hat{\sigma}_+}{dt} = -i\omega_{01}\hat{\sigma}_+ \quad (2.10)$$

And the time evolution of the operators can be written as:

$$\hat{\sigma}_-(t) = e^{-i\omega_{01}t} \hat{\sigma}_- \quad \hat{\sigma}_+(t) = e^{i\omega_{01}t} \hat{\sigma}_+ \quad (2.11)$$

Substituting this into Eq.2.7, the Hamiltonian becomes:

$$\hat{H}_d = -i\gamma \left[e^{i(\omega-\omega_{01})t} \hat{\sigma}_- - e^{i(\omega+\omega_{01})t} \hat{\sigma}_+ + e^{-i(\omega+\omega_{01})t} \hat{\sigma}_- - e^{i(-\omega+\omega_{01})t} \hat{\sigma}_+ \right] \quad (2.12)$$

Where all the real and constant terms were grouped into γ . The four terms of Eq.2.12 are oscillating at different frequencies. Terms oscillating at higher frequencies will average out faster than the timescales of the system, so they can be neglected. This is known as the *rotating-wave approximation* (RWA). Applying this, the Hamiltonian becomes:

$$\hat{H}_d = -i\gamma \left[e^{i(\omega-\omega_{01})t} \hat{\sigma}_- - e^{i(-\omega+\omega_{01})t} \hat{\sigma}_+ \right] \quad (2.13)$$

From this representation it can be seen that the perturbation term $\hat{H}_d$ contains the ladder operators $\hat{\sigma}_+$ and $\hat{\sigma}_-$, so it can induce $|0\rangle \rightarrow |1\rangle$ and $|1\rangle \rightarrow |0\rangle$ transitions. It can be proven [52] that this term is responsible for oscillations over time of the state probabilities $|\alpha|^2$ and $|\beta|^2$ introduced in Eq.1.1 (see Fig.2.1b). In other words, the microwave drive induces coherent oscillations between the ground and the excited state. These oscillations are known as *Rabi oscillations*, and are described by *Rabi's formula*. In the hypothesis that, at $t = 0$, the qubit is in the ground state, Rabi's formula is:

$$|\beta(t)|^2 = \frac{\gamma^2/\hbar^2}{\gamma^2/\hbar^2 + (\omega - \omega_{01})^2/4} \sin^2 \left\{ \left[\frac{\gamma^2}{\hbar^2} + \frac{(\omega - \omega_{01})^2}{4} \right]^{1/2} t \right\} \quad (2.14)$$

With $|\alpha(t)|^2 = 1 - |\beta(t)|^2$. The angular frequency of Rabi oscillations is:

$$\Omega = 2\sqrt{\frac{\gamma^2}{\hbar^2} + \frac{(\omega - \omega_{01})^2}{4}} \quad (2.15)$$

And it can be seen that the magnitude of the oscillation is maximum, with $|\alpha(t)|^2$ oscillating from 0 to 1, when the frequency of the driving pulse ω is exactly ω_{01} . When the drive frequency is detuned from the qubit frequency, the amplitude of the oscillations is damped following a Lorentzian shape. This can be explained by the fact that, for a detuned signal, the energy of the photons sent to the qubit is different from the transition energy of the qubit itself.

From Equation 2.14 it can be seen that, by keeping the drive signal on, $|\alpha(t)|^2$ and $|\beta(t)|^2$ oscillate indefinitely. To bring the qubit from $|0\rangle$ to $|1\rangle$, then, the drive pulse has to last only half of the Rabi period, so that $|\beta(t)|^2$ will be 1 (see Fig.2.1b). Notice that the same drive pulse, if applied when the qubit is already in $|1\rangle$, will bring its state back to $|0\rangle$. The same drive pulses can then have different outcomes depending on the qubit state. This is a general result, and it will be shown in the next section that it is a consequence of the fact that drive pulses correspond to rotations of the Bloch vector, representing the qubit state, on the Bloch sphere. In this framework, the same rotation, if applied to a different initial position, leads to a different outcome.

2.1.2 Bloch sphere representation of qubit driving

In the hypothesis of resonant driving, i.e. with $\omega = \omega_{01}$, Equation 2.13 reduces to:

$$\hat{H}_d = -i\gamma(\hat{\sigma}_- - \hat{\sigma}_+) = \gamma\hat{\sigma}_y \quad (2.16)$$

A qubit state $|\psi\rangle$ then evolves following the equation:

$$i\hbar \frac{d}{dt} |\psi(t)\rangle = \gamma \hat{\sigma}_y |\psi(t)\rangle \quad (2.17)$$

The Pauli matrices are the generators of the SU(2) group, which is a double cover of the rotation group SO(3). Eq. 2.17 is then representing a rotation of the Bloch vector associated with $|\psi\rangle$ around the y axis, induced by the microwave drive.

Drive pulses exciting the qubit from $|0\rangle$ to $|1\rangle$ or, vice versa, de-exciting the qubit from $|1\rangle$ to $|0\rangle$, are then perfectly equivalent, because they both correspond to a π -radian rotation of the Bloch vector. This is the reason why these pulses are usually called π -pulses. Following the same principle, to drive the qubit from $|0\rangle$ to a superposition state $1/\sqrt{2}|0\rangle + 1/\sqrt{2}|1\rangle$, which is very common in qubit experiments, a rotation of $\pi/2$ radians is needed. The corresponding pulse is then called $\pi/2$ -pulse. From Fig.2.1b it can be seen that a $\pi/2$ -pulse is just a π -pulse lasting half the time.

Generalizing this result, it is possible to understand what happens when sending drive pulses with different phases. Up until now, the drive pulses considered were described by Eq. 2.6. By introducing, instead, a phase ϕ :

$$V_g(t) = V_g \cos(\omega t + \phi) = \frac{V_g}{2} \left[e^{i(\omega t + \phi)} + e^{-i(\omega t + \phi)} \right] \quad (2.18)$$

Eq.2.13 becomes, again in the hypothesis of resonant driving:

$$\hat{H}_d = -i\gamma \left[e^{i\phi} \hat{\sigma}_- - e^{-i\phi} \hat{\sigma}_+ \right] = \gamma (\hat{\sigma}_x \sin \phi - \hat{\sigma}_y \cos \phi) \quad (2.19)$$

So the rotation axis is not y anymore, but it is tilted (even if still in the xy -plane). As an example, for $\phi = \pi/2$ the rotation axis becomes x , orthogonal to y .

2.1.3 Fine-tuning the drive pulse frequency

As discussed in Section 2.1.1, the amplitude of Rabi oscillations drops rapidly when the drive frequency ω gets farther from the qubit frequency ω_{01} . Fine-tuning the frequency of the drive pulse is then a fundamental step to avoid errors in computations. This is usually achieved by doing *Ramsey interferometry* [53, 54].

The principle behind this measurement is that a small frequency mismatch between the drive and the qubit leads to a relative phase that accumulates linearly over time. Ramsey

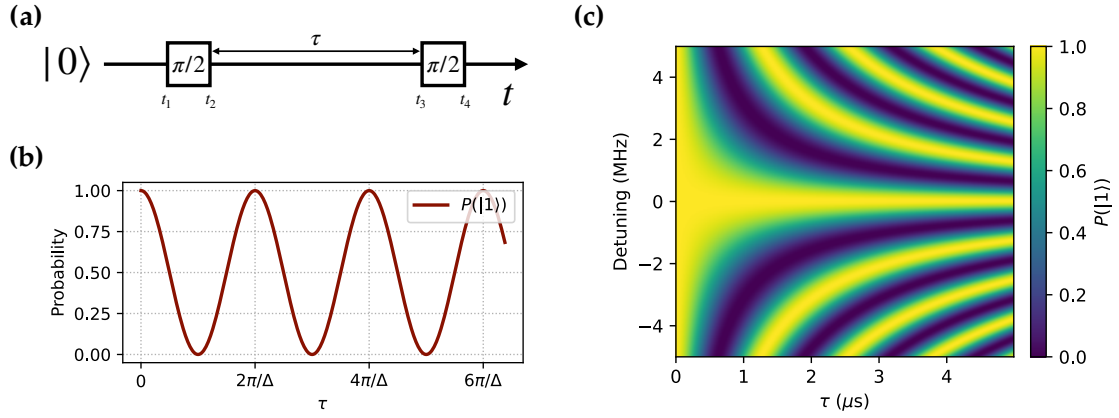


FIGURE 2.2: (a) Scheme of how a Ramsey experiment is done, (b) the oscillations of $P(|1\rangle)$ at given detuning Δ varying the delay τ and (c) simulation of $P(|1\rangle)$ when varying both the detuning and τ .

interferometry makes it possible to study the phase evolution of a qubit prepared in a superposition state and allowed to freely evolve for a certain time, thereby enabling a precise calibration of the frequency of the drive. Furthermore, it allows to monitor how coherent this evolution remains in time.

In Ramsey interferometry, two $\pi/2$ -pulses with frequency ω , separated by a time delay τ , are sent to a qubit in $|0\rangle$ (see Fig.2.2a). In the ideal case, with ω corresponding exactly to the qubit frequency ω_{01} , the two $\pi/2$ rotations will sum up and the final state of the qubit will be $|1\rangle$. If instead the drive pulse is detuned from the qubit frequency, so $\omega \neq \omega_{01}$, this is no longer valid, due to the free evolution of the qubit during τ .

To show that, four time instants during a Ramsey experiment have to be considered: t_1 , t_2 , t_3 and t_4 , which are, respectively, the time when the first $\pi/2$ -pulse starts, when it ends, when the second pulse starts and when the second pulse ends. In all the next calculations, the detuning will be considered small enough to make the effect of the two pulses still be close to a $\pi/2$ rotation.

During τ , so between t_2 and t_3 , the time evolution of the qubit state is determined by the unperturbed Hamiltonian (in the two-level approximation):

$$\hat{H}_0 = \frac{\hbar\omega_{01}}{2}\hat{\sigma}_z \quad (2.20)$$

Resulting in:

$$|\psi(t_2 + t)\rangle = \exp\left(-i\frac{\omega_{01}}{2}\hat{\sigma}_z t\right)|\psi(t_2)\rangle \quad (2.21)$$

$|\psi(t_2)\rangle$ corresponds to the qubit state at the end of the first pulse, so:

$$|\psi(t_2)\rangle = \frac{1}{\sqrt{2}}(|0\rangle + |1\rangle) \quad (2.22)$$

By applying the time-evolution operator, the qubit state at $t_3 = t_2 + \tau$ is given by:

$$|\psi(t_3)\rangle = \frac{1}{\sqrt{2}}(|0\rangle + e^{i\omega_{01}\tau}|1\rangle) \quad (2.23)$$

Where the global phase $\exp(-i\omega_{01}\tau/2)$ has been discarded. As seen earlier, a $\pi/2$ -pulse is a rotation along the y axis: this means that the second $\pi/2$ -pulse will not just send the qubit to $|1\rangle$, but the final position of the Bloch vector will be a consequence of where it will be on the Bloch sphere, so of the relative phase between the two eigenstates.

To compute this rotation it is easier to move to a reference frame rotating with the same frequency ω of the drive [55]. In this frame, $|\psi(t_3)\rangle$ becomes:

$$|\psi(t_3)\rangle = \frac{1}{\sqrt{2}}(|0\rangle + e^{i(\omega_{01}-\omega)\tau}|1\rangle) = \frac{1}{\sqrt{2}}(|0\rangle + e^{i\Delta\tau}|1\rangle) \quad (2.24)$$

Where $\Delta = \omega_{01} - \omega$ is the detuning between the qubit transition frequency and the drive frequency. If $\Delta = 0$, so the applied pulses are at resonance with the qubit, the Bloch vector does not rotate during τ , and the second $\pi/2$ -pulse will just complete the π rotation around y , bringing the qubit to $|1\rangle$.

If $\Delta \neq 0$, the $\pi/2$ rotation around y is applied to the state $|\psi(t_3)\rangle$ written in Eq. 2.24. The rotation operator is, in matrix form:

$$R_y(\pi/2) = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & -1 \\ 1 & 1 \end{pmatrix} \quad (2.25)$$

And, by applying it to $|\psi(t_3)\rangle$, the final state is computed:

$$|\psi(t_4)\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & -1 \\ 1 & 1 \end{pmatrix} \frac{1}{\sqrt{2}} \begin{pmatrix} 1 \\ e^{i\Delta\tau} \end{pmatrix} = \frac{1}{2} \begin{pmatrix} 1 - e^{i\Delta\tau} \\ 1 + e^{i\Delta\tau} \end{pmatrix} \quad (2.26)$$

By factorizing and removing a global phase term $e^{i\Delta\tau/2}$, the state becomes:

$$|\psi(t_4)\rangle = -i \sin\left(\frac{\Delta\tau}{2}\right)|0\rangle + \cos\left(\frac{\Delta\tau}{2}\right)|1\rangle \quad (2.27)$$

From which the measurement probabilities can be computed:

$$P(|1\rangle) = \cos^2\left(\frac{\Delta \cdot \tau}{2}\right) \quad P(|0\rangle) = \sin^2\left(\frac{\Delta \cdot \tau}{2}\right) \quad (2.28)$$

So, when the qubit is detuned, the probability to have the qubit in $|1\rangle$ after the second $\pi/2$ pulse oscillates with frequency $\Delta/4$ while varying τ (see Fig.2.2b). The probability is 1 when $\tau = 2k\pi/\Delta$, with k positive integer, and is 0 when $\tau = (2k+1)\pi/\Delta$. On the Bloch sphere, the first condition corresponds to when the Bloch vector has made multiple full precessions and is back in $1/\sqrt{2}(|0\rangle + |1\rangle)$. When the second condition is true, instead, the Bloch vector has undergone an additional half precession and is in $1/\sqrt{2}(|0\rangle - |1\rangle)$, so a $\pi/2$ rotation around y brings it back to $|0\rangle$.

At the same time, $P(|1\rangle)$ oscillates also when keeping τ fixed and changing the detuning Δ . In this case the oscillation frequency is $\tau/4$, so the width of the fringes can be varied by changing the delay τ . Fig.2.2c shows the Ramsey fringes obtained when varying both Δ and τ .

Since the central fringe is always centered at the qubit frequency ($\Delta = 0$), but its width can be narrowed by varying τ , this method allows a very precise estimation of ω_{01} . Because of that, for all the qubits studied in this thesis project, Ramsey interferometry has been a fundamental step in their characterization.

Summarizing, Rabi oscillation measurements (Section 2.1.1) are used to calibrate the amplitude and duration of control pulses, ensuring accurate single-qubit rotations. Ramsey interferometry, on the other hand, probes the phase evolution of the qubit under free evolution, providing a precise determination of its resonance frequency.

2.2 Qubit state readout

According to the laws of quantum mechanics, every measurement on a quantum system inevitably disturbs the system itself [52]. Measuring the state of a quantum system makes the wavefunction collapse into one of its eigenstates, so any information regarding superposition states can only be inferred statistically. At the same time, the uncertainty relation states that, if the system is characterized by non-commuting variables, increasing the precision of the measurement on one variable will automatically result in a higher uncertainty on the other. In the worst-case scenario, the measurement process disturbs the quantum system so much that all subsequent measurements will be affected. This is a typical problem of quantum optics experiments, where measurements are often done with photocounting techniques that destroy the quantum state [56].

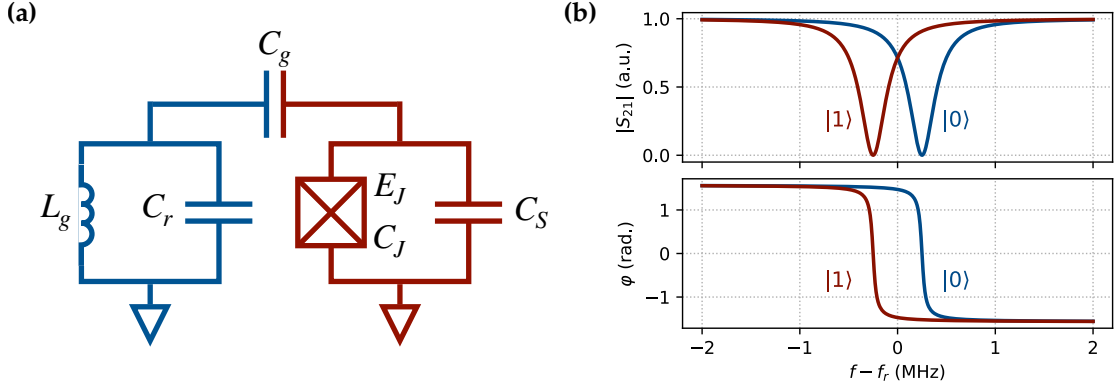


FIGURE 2.3: (a) Scheme of a transmon qubit capacitively coupled to a LC resonator and (b) dispersive shift in amplitude (top) and phase (bottom).

In this scenario, measurements on superconducting qubits are no exception. For reliable quantum computations, then, it is crucial to employ a measurement technique that does not affect the outcome of the operations that will be performed afterwards. In other words, the outcome of a qubit measurement must be a consequence solely of the quantum state being measured and not of the measurement process. Techniques satisfying this requirement are known as *quantum nondemolition measurements* or QND measurements [57, 58].

The most common measurement technique for superconducting qubits is *dispersive shift readout* [59], where the qubit is coupled to a resonator and induces a state-dependent shift in its resonance frequency. By measuring this frequency shift, one can infer the quantum state without destroying it. Alternatives to this approach have been proposed [60, 61], but dispersive readout remains the most common and experimentally reliable technique.

2.2.1 Dispersive shift readout

To show the working principle of dispersive shift, a transmon qubit capacitively coupled to a planar LC resonator (see Fig. 2.3a) will be considered. Notice, however, that this type of readout also works with other superconducting qubit designs [37, 47, 49], different types of resonators, like 3D cavities [62], and different couplings, like shared inductances [63].

The Hamiltonian for the circuit in Fig.2.3a. can be derived in a similar way to what has been done for qubit driving [21]. While for qubit driving the generator was introducing only a static offset term, though, when the qubit is coupled to a resonator a new charge term appears, this time acting as a dynamical variable Q_r . Before quantization the Hamiltonian is:

$$H = \frac{1}{2C_\Sigma} \left(Q + \frac{C_g}{C_r} Q_r \right)^2 - E_J \cos \varphi + \frac{Q_r^2}{2C_r} + \frac{\Phi_r}{2L_r} \quad (2.29)$$

Where C_r and L_r are, respectively, the capacitance and the inductance of the resonator, C_g is the coupling capacitance, C_J the intrinsic capacitance of the junction, C_S the shunting capacitance parallel to the junction and $C_\Sigma = C_S + C_J$. This Hamiltonian can be rearranged as:

$$H = \underbrace{\frac{Q^2}{2C_\Sigma} - E_J \cos \varphi}_{\text{qubit}} + \underbrace{\left(\frac{1}{2C_\Sigma} \frac{C_g^2}{C_r^2} + \frac{1}{2C_r} \right) Q_r^2 + \frac{\Phi_r}{2L_r}}_{\text{resonator}} + \underbrace{\frac{C_g}{C_\Sigma C_r} Q Q_r}_{\text{coupling}} \quad (2.30)$$

Resulting in three terms: a qubit term, a resonator term (with different charging energy due to the coupling to the qubit) and a coupling term proportional to C_g . Like in Eq.1.26, the cosine term can be Taylor expanded up to fourth order:

$$H \approx \frac{Q^2}{2C_\Sigma} + \frac{1}{2} E_J \varphi^2 - \frac{1}{24} E_J \varphi^4 + \left(\frac{1}{2C_\Sigma} \frac{C_g^2}{C_r^2} + \frac{1}{2C_r} \right) Q_r^2 + \frac{\Phi_r}{2L_r} + \frac{C_g}{C_\Sigma C_r} Q Q_r \quad (2.31)$$

And the Hamiltonian can be quantized by promoting the variables to operators. Operators can then be rewritten in terms of the ladder operators introduced in sections 1.2.1 and 1.2.2, calling $\hat{a}$ and $\hat{a}^\dagger$ the operators of the resonator and $\hat{b}$ and $\hat{b}^\dagger$ the operators of the qubit. In order to simplify the φ^4 term, the *rotating wave approximation* introduced in Section 2.1.1 is applied again, allowing to keep only the terms with equal number of $\hat{b}$ and $\hat{b}^\dagger$. Under these assumptions, the Hamiltonian becomes¹:

$$\hat{H} \approx \underbrace{\hbar\omega_{01}\hat{b}^\dagger\hat{b} - \frac{E_C}{2}\hat{b}^\dagger\hat{b}^\dagger\hat{b}\hat{b}}_{\text{qubit}} + \underbrace{\hbar\omega_r\hat{a}^\dagger\hat{a}}_{\text{resonator}} - \underbrace{\hbar g(\hat{b}^\dagger - \hat{b})(\hat{a}^\dagger - \hat{a})}_{\text{coupling}} \quad (2.32)$$

The rotating wave approximation can be applied also to the coupling term, which simplifies as $\approx \hbar g(\hat{b}^\dagger\hat{a} + \hat{b}\hat{a}^\dagger)$. By representing the qubit as a two-level system, like in Section 2.1.1, $\hat{b}$ and $\hat{b}^\dagger$ can be replaced with Pauli matrices, obtaining the *Jaynes-Cummings Hamiltonian* [64]:

$$\hat{H}_{JC} = \frac{\hbar\omega_{01}}{2}\hat{\sigma}_z + \hbar\omega_r\hat{a}^\dagger\hat{a} + \hbar g(\hat{a}^\dagger\hat{\sigma}_- + \hat{a}\hat{\sigma}_+) \quad (2.33)$$

¹To be rigorous, every resonator is resonant not only at ω_r but also at its integer multiples, so higher modes should also be taken into account. In most qubit setups, though, the qubit frequency is much closer to the first mode than to the others, so the contribution of higher modes will be much smaller and can be safely neglected.

In this Hamiltonian the coupling term describes quanta exchanges between the qubit and the resonator. When qubit and resonator are at resonance, i.e. when $\Delta = \omega_{01} - \omega_r = 0$, this corresponds to an energy exchange that disrupts the qubit state, which is the opposite of what a QND measurement requires. It is far more convenient, then, to work in the *dispersive regime*, where the detuning Δ is much larger than the coupling strength g . Here, the interaction can take place only through virtual photons, and the effect of the coupling term can be estimated with a perturbative approach [30]. In this regime the Hamiltonian becomes:

$$\hat{H} = \hbar(\omega_r + \chi\hat{\sigma}_z)\left(\hat{a}^\dagger\hat{a} + \frac{1}{2}\right) + \frac{\hbar\tilde{\omega}_{01}}{2}\hat{\sigma}_z \quad (2.34)$$

Where $\chi = g^2/\Delta$ and $\tilde{\omega}_{01} = \omega_{01} + \chi$. Here, χ plays two roles: it produces a *Lamb shift* in the qubit frequency and, more important, it introduces a shift of the resonator frequency whose sign is determined by the eigenvalue of $\hat{\sigma}_z$, so by the qubit state (see Fig.2.3b). The qubit state can then be measured, without being demolished, by probing the resonator at a specific frequency and looking at the transmitted signal, which value will depend on the position of the resonance.

2.2.2 Signal demodulation and readout

Now that a method for measuring the qubit state without demolishing it has been established, the next step is to perform the measurement experimentally. To probe the resonator, a microwave pulse is sent to the chip. The frequency ω_{RO} of the readout signal is chosen in such a way that the discrimination between $|0\rangle$ and $|1\rangle$ is maximum (see next section). Following [22], the transmitted signal can be written as:

$$s(t) = A_{RO} \cos(\omega_{RO}t + \theta_{RO}) \quad (2.35)$$

Since ω_{RO} is fixed, all the information about the qubit state is stored in the amplitude A_{RO} and in the phase θ_{RO} (see Fig.2.3b). By rewriting the signal as

$$s(t) = \Re\left\{A_{RO}e^{i(\omega_{RO}t + \theta_{RO})}\right\} = \Re\left\{A_{RO}e^{i\theta_{RO}}e^{i\omega_{RO}t}\right\} \quad (2.36)$$

The quantity of interest in this analysis is the phasor $A_{RO}e^{i\theta_{RO}}$. This can be decomposed into a real and an imaginary part, called respectively the *in-phase* I and the *quadrature* Q :

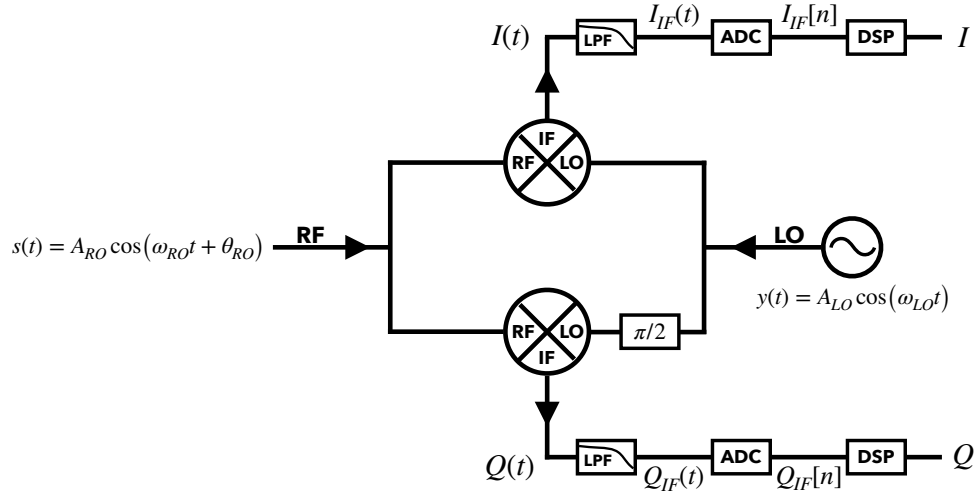


FIGURE 2.4: Scheme of the demodulation and readout of a RF signal.

$$A_{RO}e^{i\theta_{RO}} = A_{RO} \cos \theta_{RO} + iA_{RO} \sin \theta_{RO} = I + iQ \quad (2.37)$$

A typical readout setup for measuring I and Q is depicted in Fig.2.4. The incoming signal is split in two and sent to two frequency mixers, where it gets mixed with the signal coming from a local oscillator (LO). Defining the LO signal as:

$$y(t) = A_{LO} \cos(\omega_{LO}t) \quad (2.38)$$

The signal coming out of the mixer on the I branch is:

$$I(t) = \frac{1}{4}s(t)y(t) = \frac{A_{RO}A_{LO}}{4} \cos(\omega_{RO}t + \theta_{RO}) \cos(\omega_{LO}t) \quad (2.39)$$

Which yields:

$$\begin{aligned} I(t) = \frac{A_{RO}A_{LO}}{8} \cos \theta_{RO} & \left\{ \cos[(\omega_{LO} + \omega_{RO})t] + \cos[(\omega_{RO} - \omega_{LO})t] \right\} \\ & - \frac{A_{RO}A_{LO}}{8} \sin \theta_{RO} \left\{ \sin[(\omega_{RO} + \omega_{LO})t] + \sin[(\omega_{RO} - \omega_{LO})t] \right\} \end{aligned} \quad (2.40)$$

Something similar is done also on the Q branch, but there the LO signal is phase-shifted by $\pi/2$ before being mixed, resulting in:

$$Q(t) = \frac{A_{RO}}{2} \cos(\omega_{RO}t + \theta_{RO}) \cdot \left(-\frac{A_{LO}}{2}\right) \sin(\omega_{LO}t) \quad (2.41)$$

This yields:

$$Q(t) = -\frac{A_{RO}A_{LO}}{8} \cos \theta_{RO} \left\{ \sin[(\omega_{RO} + \omega_{LO})t] + \sin[(\omega_{RO} - \omega_{LO})t] \right\} \\ + \frac{A_{RO}A_{LO}}{8} \sin \theta_{RO} \left\{ \cos[(\omega_{RO} + \omega_{LO})t] + \cos[(\omega_{RO} - \omega_{LO})t] \right\} \quad (2.42)$$

The obtained $I(t)$ and $Q(t)$ have two frequency components, one with $\omega_{RO} + \omega_{LO}$ and the other with $\omega_{RO} - \omega_{LO}$. The high-frequency component gets filtered out with a low-pass filter, while the low-frequency one gets digitized and processed.

Depending on the choice of ω_{LO} there are two different strategies for getting I and Q : *homodyne demodulation* and *heterodyne demodulation*.

Homodyne demodulation is the simplest one: ω_{LO} is chosen to be equal to ω_{RO} , making the oscillating terms vanish and leaving only:

$$I = \frac{A_{RO}A_{LO}}{8} \cos \theta_{RO} \quad Q = \frac{A_{RO}A_{LO}}{8} \sin \theta_{RO} \quad (2.43)$$

From which amplitude and phase of the signal can be computed:

$$A_{RO} \propto \sqrt{I^2 + Q^2} \quad \theta_{RO} = \arctan(Q/I) \quad (2.44)$$

Despite the simplicity in getting A_{RO} and θ_{RO} , homodyne demodulation has two major drawbacks that makes its choice unfavorable. The first one is that is highly affected by *Flicker noise* or *1/f noise*: it is a type of electronic noise typical of many devices such as solid-state components and vacuum tubes, characterized by a power spectral density going as the inverse of the frequency [65]. The second drawback is that this type of demodulation does not work with *frequency division multiplexing* (FDM), i.e. when a single pulse is composed by different carrier frequencies:

$$s(t) = \sum_{i=1}^N A_{RO}^{(i)} \cos[\omega_{RO}^{(i)}t + \theta^{(i)}] \quad (2.45)$$

Frequency multiplexing is very useful in quantum computing, because it allows to measure simultaneously the state of multiple qubits, so not being able to demodulate these signals is a huge limitation for qubit experiments.

Both these issues are addressed with heterodyne demodulation, where $\omega_{LO} \neq \omega_{RO}$. Here the oscillating terms in Eq.2.40 and Eq.2.42 do not vanish completely, so when filtering out the high-frequency components the result is:

$$I_{IF}(t) = \frac{A_{RO}A_{LO}}{8} \cos(\omega_{IF}t + \theta_{RO}) \quad (2.46)$$

$$Q_{IF}(t) = \frac{A_{RO}A_{LO}}{8} \sin(\omega_{IF}t + \theta_{RO}) \quad (2.47)$$

Where $\omega_{IF} = |\omega_{RO} - \omega_{LO}|$ is called the *intermediate frequency*. Now the two terms are still oscillating, but ω_{IF} is small enough to allow digitization. During digitization the signals are sampled with a time interval Δt between consecutive samples. The result can be rewritten in complex form as:

$$z_{IF}[n] = I_{IF}[n] + iQ_{IF}[n] = \frac{A_{RO}A_{LO}}{8} \left[\cos(\Omega_{IF}n + \theta_{RO}) + i \sin(\Omega_{IF}n + \theta_{RO}) \right] \quad (2.48)$$

Where $n = t/\Delta t$ is the number of samples and $\Omega_{IF} = \omega_{IF}\Delta t$ is the digital frequency. $z_{IF}[n]$ can be written equivalently as:

$$z_{IF}[n] = \frac{A_{RO}A_{LO}}{8} e^{i\theta_{RO}} e^{i\Omega_{IF}n} \quad (2.49)$$

$\exp(i\Omega_{IF}n)$ can be removed with a point-by-point multiplication:

$$z[n] = z_{IF}[n] * e^{-i\Omega_{IF}n} \quad (2.50)$$

And by averaging over all the points the outcome is:

$$\bar{z}[n] = \frac{1}{N} \sum_n z[n] = \frac{A_{RO}A_{LO}}{8} e^{i\theta_{RO}} \quad (2.51)$$

From which I and Q can be related to A_{RO} and θ_{RO} :

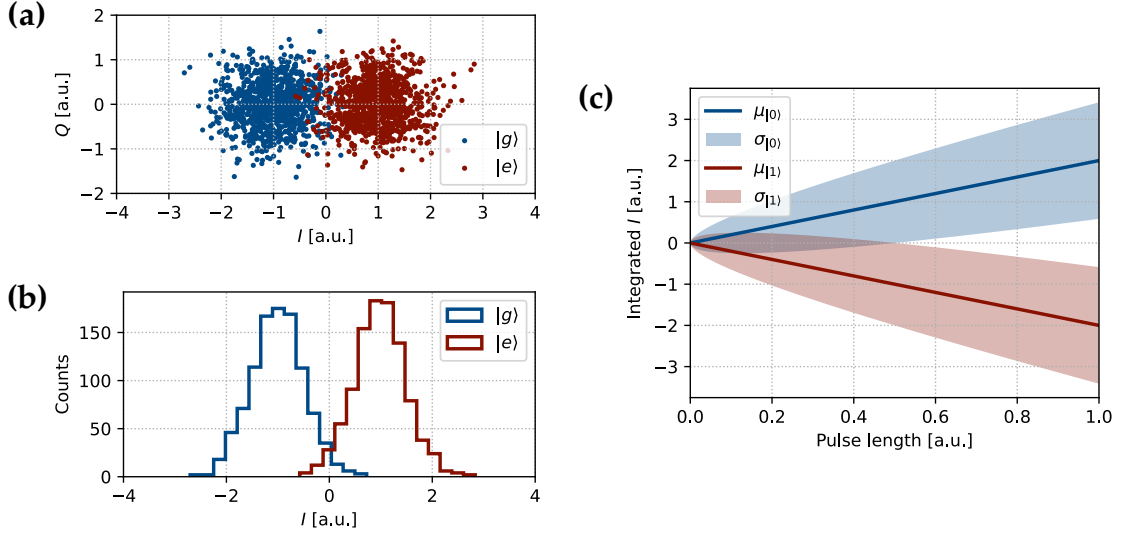


FIGURE 2.5: Simulation of the dispersive shift measurement of a qubit state: (a) clouds on the IQ plane produced by repeatedly measuring the qubit when in $|g\rangle$ and when in $|e\rangle$, (b) the histograms of the I values of the two clouds, and (c) mean I value and its standard deviation as a function of the pulse length for both states.

$$I = \frac{A_{RO}A_{LO}}{8} \cos \theta_{RO} \quad Q = \frac{A_{RO}A_{LO}}{8} \sin \theta_{RO} \quad (2.52)$$

A_{RO} and θ_{RO} are then obtained like in Eq.2.44.

With heterodyne demodulation Flicker noise is much damped since I_{IF} and Q_{IF} are oscillating and not DC. At the same time, multiplexed signals can now be demodulated since each tone will have a different $\omega_{IF}^{(i)}$, and the tones can be isolated with band-pass filters and digitized individually.

2.2.3 Readout fidelity, noise and signal amplification

Every measurement, including qubit state readout, is affected by noise. In the absence of noise, the only requirement for distinguishing $|0\rangle$ and $|1\rangle$ would be to tune the readout pulse at a frequency where the response of the circuit is state-dependent. By repeatedly measuring the qubit state and then plotting the response in a plane with I and Q as its axes (which, as seen in the last section, is equivalent to looking at A_{RO} and θ_{RO}), ideally only two dots would appear, one when the qubit is in $|0\rangle$ and the other when the qubit is in $|1\rangle$. This is not what happens in real experiments, though: values of I and Q fluctuate due to various sources of noise, and an ensemble of measurements would instead produce two gaussian distributions corresponding to the two qubit states, as can be seen in Fig.2.5a and Fig.2.5b. If these two gaussians partially overlap, state discrimination may not be always correct.

The probability to correctly identify the qubit state in a readout measurement is called *assignment fidelity* and, in the hypothesis that the qubit was prepared 50% of the times in $|0\rangle$ and 50% of the times in $|1\rangle$, is defined as [66]:

$$\mathcal{F} = 1 - \frac{1}{2} [P(1|0) + P(0|1)] \quad (2.53)$$

Where $P(1|0)$ is the probability to measure the qubit in $|e\rangle$ when it was actually in $|g\rangle$, and $P(0|1)$ is, vice versa, the probability to measure the qubit in $|g\rangle$ when it was actually in $|e\rangle$. In some works, the *measurement fidelity* is used instead of the assignment fidelity. Measurement fidelity is substantially the same as assignment fidelity, but the factor $1/2$ is missing [67]. Due to its definition, measurement fidelity ranges from -1 to +1, while assignment fidelity, being a probability, ranges from 0 to +1.

Achieving a high fidelity is fundamental to avoid computational errors, and as of today experiments have achieved single-shot readout fidelities of up to approximately 99.8% under optimized laboratory conditions [68–71].

As can be seen from Fig.2.5, the readout fidelity is a consequence of the distance between the centers of the two distributions, $\mu_{|0\rangle}$ and $\mu_{|1\rangle}$, and of their width, i.e. their standard deviations, $\sigma_{|0\rangle}$ and $\sigma_{|1\rangle}$. These three variables determine the *signal-to-noise ratio* (SNR), which is defined as [22]:

$$\text{SNR} = \frac{|\mu_{|0\rangle} - \mu_{|1\rangle}|}{2(\sigma_{|0\rangle} + \sigma_{|1\rangle})} \quad (2.54)$$

Improving the signal-to-noise ratio is an effort that takes into account different aspects of a qubit experiment, from the qubit-resonator coupling strength to the optimization of the readout pulse and the amplification of the returning signal. The qubit-resonator coupling is a design choice and so it cannot be modified once the device has been produced, but pulse optimization and signal amplification can be done whatever is the setup studied.

Concerning the readout pulse optimization, the variables playing a role in the fidelity are the pulse frequency, the pulse amplitude and the pulse length. Pulse frequency has to be chosen in such a way that $|\mu_{|0\rangle} - \mu_{|1\rangle}|$ is maximum. The best value depends on the variable chosen for state discrimination: as shown in Fig.2.3, when the phase is measured the maximum separation between the qubit states occurs half-way between the qubit-dependent resonant frequencies, so at $\omega_{RO} = (\omega_r^{(0)} + \omega_r^{(1)})/2$; when measuring the amplitude instead, the biggest separation between the two states occurs at $\omega_{RO} = \omega_r^{(0)}$ or $\omega_{RO} = \omega_r^{(1)}$.

Variations in the pulse amplitude and in the pulse length have a similar effect on the fidelity. This behavior can be intuitively understood from Eq.2.34, where the dispersive shift term is given by:

$$\hat{H}_{disp} = \hbar\chi\hat{\sigma}_z\hat{a}^\dagger\hat{a} \quad (2.55)$$

So the more photons are sent to probe the resonator, the higher will be the contribution of $\hat{H}_{disp}$, since $\hat{a}^\dagger\hat{a}$ is the number operator, the larger will be the separation $|\mu_{|0\rangle} - \mu_{|1\rangle}|$ between the two clouds. At the same time, fluctuations in the number of photons (what is usually called *shot noise*) will follow a Poisson distribution, so $\sigma_{|0\rangle}$ and $\sigma_{|1\rangle}$ will increase slower, as the square root of the total number [72].

Increasing the pulse amplitude or the pulse length therefore improves the separation between the two clouds and, consequently, the signal-to-noise ratio (see Fig.2.5c). However, these variables cannot be increased indefinitely, for two main reasons. The first one is that, if the probe is too strong, the qubit-induced frequency shift becomes negligible compared to the resonator linewidth (see eq.2.33), effectively decoupling the qubit from the cavity. In this so-called *punch out* regime [21], the resonator response no longer carries information about the qubit state. At the same time, keeping the pulse amplitude low while increasing the pulse length too much is also problematic: a reliable readout must occur on a timescale much shorter than the characteristic time of the free evolution of the system. This means that pulses have to be short in time [73], with typical values for superconducting qubits ranging from tens of ns up to μs .

The last variable to be considered for reaching high fidelities is signal amplification. This is a fundamental step in qubit state readout, because the signals probing resonator are extremely weak, typically around $-110 \div -100$ dBm for μs -long pulses. Trying to demodulate a signal so small with the electronics used in qubit experiments would result in observing only noise. This is due to the fact that the few photons of a readout signals would be completely drowned in the blackbody radiation of the readout electronics [74]. This radiation is produced by the thermal vibrations of charges in the setup, and is known as *Johnson noise* or *thermal noise* [65].

The signal needs then to be amplified before being demodulated, but this cannot be done at room temperature, because again the thermal noise added by the amplifier would overwhelm the returning pulse. The solution is to amplify already at cryogenic temperatures (see Fig. 2.6), where the thermal noise is much smaller. Its power spectral density, in fact, goes as:

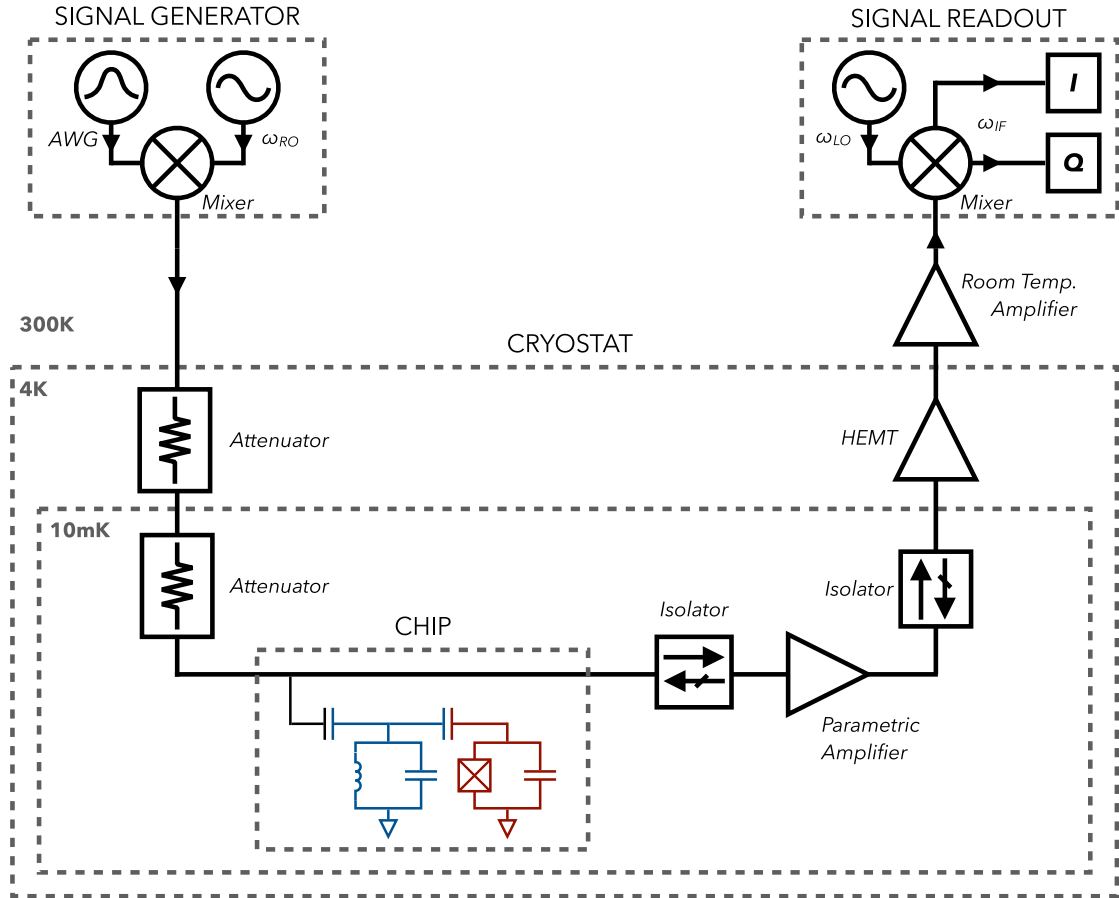


FIGURE 2.6: Simplified scheme of the experimental setup of a qubit experiment. Notice that a lot of elements are missing from this representation, including filters, DC lines for flux biasing and the RF line for the pump tone of the parametric amplifier.

$$S(T, \omega) \propto \frac{\hbar\omega}{e^{\hbar\omega/k_B T} - 1} \quad (2.56)$$

So, for $\omega = 5$ GHz, the spectral density at 20 mK is five orders of magnitude lower than at room temperature².

Amplifying at cryogenic temperatures, though, is not simple, because the typical cooling power of a dilution refrigerator at 20 mK is lower than $10 \mu\text{W}$ [25], while most of RF amplifiers commercially available dissipate much more than that. Also, every amplifier must be well isolated from the chip under test, in order to prevent the noise that it generates from reaching the qubit. As of today, the state-of-the-art cryogenic amplifiers are *superconducting parametric amplifiers* [75], parametric amplifiers where the incoming signal is amplified by mixing it with a strong tone, usually called *pump*. The idea is that,

²This is the same reason why producing a pulse so weak at room temperature would also result in observing just noise after demodulation. The solution for that is to generate a much stronger pulse at room temperature and attenuate it as the temperature drops, in order to remove the blackbody radiation of higher temperature stages (see Fig. 2.6)

in the amplifier, energy is drained from the pump and converted into signal photons. This is done in several ways, but the general idea is similar for all the designs.

To understand the working principle of superconducting parametric amplifiers, a good starting point is a LC resonator coupled to a transmission. A system like that does not do much to incoming signals, except for introducing a delay in signals with frequency close to $f_0 = 1/\sqrt{LC}$. If the inductance L is modulated as $L = L_0 + \delta L \cos(2\pi f_p t + \vartheta_p)$, though, the incoming signal can actually get amplified, provided that $f_p \approx 2f_0$ [76]. The most common way to get a variable inductance is to replace the inductor with a Josephson Junction, resulting in a design identical to a qubit. The modulation is provided by biasing the junction with a RF signal named *pump tone*. It is also common to use a SQUID instead of a single junction in order to make it tunable with a DC bias (otherwise the central frequency of the gain curve would be fixed) [77]. This type of amplifier is known as *Josephson Parametric Amplifier* (JPA).

JPA's have the advantage of adding very low noise to the signal, close to the *quantum limit* (i.e. noise arising from the quantization of photons and charges and the Heisenberg's uncertainty principle) [78–80], while still providing typical gains of approx. 20 dB. Their bandwidth, though, is very narrow, typically tens of MHz, so they are unsuitable for amplifying multiplexed signals from different qubits, whose readout frequencies can be separated also by hundreds of MHz.

A common alternative to JPAs when working with multiplexed signals is the *Traveling Wave Parametric Amplifier* (TWPA) [81–85]. TWPAs grant similar gains but over a GHz-wide bandwidth, paying the price of a slightly worse noise performance [86]. Typical designs consist in arrays of Josephson junctions and, instead of working in reflection like JPAs, they work in transmission, following a principle similar to the one of optical parametric amplifiers [87, 88].

Amplification at the coldest level with JPAs and TWPAs is not enough, though, to have a good SNR, so the signal is amplified also at higher temperature stages. Typical amplification stages are at 4K with a *High Electron Mobility Transistor* (HEMT) amplifier [65] and, finally, at room temperature before the readout.

2.3 Decoherence

Every algorithm running on quantum computers consists in preparing a set of N qubits into some states $\{|\psi\rangle_n\}_{n=1,\dots,N}$, manipulate them with a sequence of unitary gates U_i

and, then, measure their final state. The success of the computation relies on the assumption that only the gates applied by the user affect the qubit states, which otherwise would stay unchanged.

However, no physical qubit that has ever been realized is capable of maintaining its state indefinitely. Superconducting qubits, when prepared into a state $|\psi\rangle$, are able to keep it only for tens or hundreds of microseconds [89], before evolving unpredictably and losing the information they were storing. Only recently some setups, specifically designed for this purpose, reached the millisecond scale [90], but even the best performing prototypes are not ready yet for most of quantum computing applications. This does not mean that research got into a dead end: to put things in perspective, in the first Cooper-pair boxes this characteristic time was $O(\text{ns})$ [28], meaning that, over the past 25 years, this parameter has improved by approximately six orders of magnitude.

The reason for this limitation lies in the interaction of qubits with the surrounding environment. In an ideal scenario qubits should be sensitive only to the signals sent by the users to manipulate and measure their state, but in reality their state is very fragile and can be affected by various factors [91], ranging from microscopic defects and oxides on the chip surface to broken Cooper pairs produced by infrared photons or environmental radioactivity. This phenomenon is known as *decoherence*, and the average time the qubit is capable of maintaining its state is known as *coherence time*. As will be discussed in the next sections, different characteristic times for qubit coherence are considered depending on the decoherence pathway.

Given the fact that decoherence affects every quantum processor, two complementary strategies have been adopted to minimize its impact:

- implement error correction algorithms into quantum computations;
- engineer qubits in order to minimize their susceptibility to external noise and maximize their coherence time.

Quantum error correction is likely to become a fundamental component of every future quantum processor, but to properly work some requirements must be satisfied. In particular, the error rate of physical qubits must be lower than a threshold [92], whose value varies with the correction algorithm considered, and errors have to be uncorrelated [93]. In order to implement error correction codes, then, qubits must have coherence times long enough to reduce the error rate below the threshold, and have to be protected against sources of correlated errors, including environmental radioactivity (see Section 3.2).

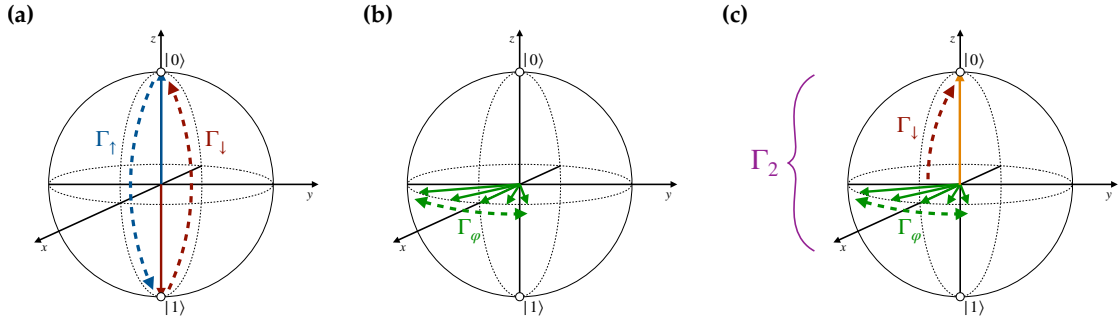


FIGURE 2.7: Decoherence in qubits: (a) longitudinal relaxation, (b) pure dephasing and (c) transverse relaxation.

2.3.1 Energy relaxation

When a qubit exchanges energy with the environment, the exchange can result in a transition from $|0\rangle$ to $|1\rangle$ or, vice versa, from $|1\rangle$ to $|0\rangle$, depending on whether the qubit is emitting a quanta or absorbing one. This type of error is known as *energy relaxation* or *longitudinal relaxation*, because on the Bloch sphere it corresponds to a depolarization along the longitudinal axis (see Fig.2.7a).

Naming $\Gamma_{\uparrow}$ the rate of $|0\rangle \rightarrow |1\rangle$ transitions and $\Gamma_{\downarrow}$ the rate of $|1\rangle \rightarrow |0\rangle$ transitions, the *longitudinal relaxation rate* Γ_1 is defined as:

$$\Gamma_1 = \frac{1}{T_1} = \Gamma_{\uparrow} + \Gamma_{\downarrow} \quad (2.57)$$

The inverse of Γ_1 is the *energy-relaxation time*, or T_1 , and corresponds to the average time the qubit spends before incurring energy relaxation. This is the same T_1 described in 1.2.4 when discussing heavy fluxonium qubits.

In the case of superconducting qubits, $|1\rangle \rightarrow |0\rangle$ transitions are much more frequent than $|0\rangle \rightarrow |1\rangle$ transitions. This is due to the fact that these qubits are operated inside dilution refrigerators with temperatures $T \approx 20$ mK. At thermal equilibrium, the transition rates are governed by Boltzmann statistics, so:

$$\frac{\Gamma_{\uparrow}}{\Gamma_{\downarrow}} = \exp(-\hbar\omega_{01}/k_bT) \quad (2.58)$$

And for $T \approx 20$ mK and $\omega_{01} \approx 5$ GHz the ratio is $\Gamma_{\uparrow}/\Gamma_{\downarrow} < 10^{-5}$. Because of that, energy relaxation is often referred also as *energy decay*.

2.3.2 Transverse relaxation

Up until now, a generic qubit state $|\psi\rangle$ has been described solely by a polar angle ϑ and an azimuthal angle φ :

$$|\psi\rangle = \cos \frac{\vartheta}{2} |0\rangle + e^{i\varphi} \sin \frac{\vartheta}{2} |1\rangle \quad (2.59)$$

With the relative phase being determined by φ . However, if the $|0\rangle$ and $|1\rangle$ states have different energies, as in superconducting qubits, there is an extra phase precession to be taken into account, already encountered when discussing Ramsey interferometry (see Section 2.1.3). In this case, in fact, the Bloch vector precesses around the z axis with frequency ω_{01} , and its correct representation is:

$$|\psi(t)\rangle = \cos \frac{\vartheta}{2} |0\rangle + e^{i\omega_{01}t} e^{i\varphi} \sin \frac{\vartheta}{2} |1\rangle \quad (2.60)$$

The Bloch vector is rarely represented in this form, because usually it is preferred to choose a frame which is rotating around the z axis at the same frequency of the qubit, but is useful to understand what happens when external noise affects the transition frequency ω_{01} .

As a matter of facts, when ω_{01} is affected by noise, its fluctuations will result in an extra phase $\exp(i\delta\omega t)$, and since these fluctuations are uncontrolled, the information about φ will be lost. This is known as *pure dephasing* (see Fig.2.7b).

The combination of pure dephasing and energy relaxation is known as *transverse relaxation* (Fig.2.7c).

As for energy relaxation, also for transverse relaxation a relaxation rate Γ_2 can be defined, which is a combination of the longitudinal relaxation rate Γ_1 and of pure dephasing rate Γ_φ :

$$\Gamma_2 = \frac{1}{T_2} = \frac{\Gamma_1}{2} + \Gamma_\varphi \quad (2.61)$$

2.3.3 Qubit relaxation measurements

Having defined the relaxation pathways for qubits, the next step is defining a measurement strategy for the characteristic times T_1 and T_2 .

Concerning T_1 , it is the sum of two terms, one for $|1\rangle \rightarrow |0\rangle$ transitions and the other for $|0\rangle \rightarrow |1\rangle$. At cryogenic temperature, though, the rate of $|0\rangle \rightarrow |1\rangle$ transitions is much lower than the other, so it can be safely neglected. The starting point for the measurement of the decay time is then a set of N identical qubits prepared in $|1\rangle$ at time $t = 0$. As time goes by, some of them will spontaneously decay to $|0\rangle$, and the variation over time of the number of qubits in $|1\rangle$ will be governed by $\Gamma_{\downarrow} \approx 1/T_1$:

$$\frac{dN}{dt} = -\Gamma_{\downarrow} N \quad (2.62)$$

Which leads to the equation:

$$N(t) = N(0)e^{-t/T_1} \quad (2.63)$$

So, by preparing a qubit in $|1\rangle$, waiting some time τ and then measuring its state, the probability to find the qubit still in $|1\rangle$ goes as:

$$P_{|1\rangle}(t) = \frac{N(t)}{N(0)} = e^{-\tau/T_1} \quad (2.64)$$

The typical way to measure T_1 consists in reconstructing the exponential probability: the qubit is prepared in $|1\rangle$, but then its state is measured only after a delay τ . By repeating the measurement several times it is possible to estimate $P_{|1\rangle}(\tau)$, and then, by repeating this procedure at different τ values, the exponential decay of $P_{|1\rangle}$ can be measured (see Fig.2.8a).

This method is very simple, but has some drawbacks: first of all, it neglects $\Gamma_{\uparrow}$, so to measure both transition rates a different strategy has to be used; other than that, since for every value of τ the measurement has to be repeated hundreds of times, the overall process takes usually tens of seconds, ending with a result which is insensitive to fast T_1 fluctuations. This is the reason why some experiments have developed different approaches, and a couple of them will be illustrated in the next chapters.

Moving on to T_2 , as seen in the last section, its value is related not only to the energy-relaxation rate Γ_1 , but also to the dephasing rate Γ_{φ} . In order to measure it, then, it is necessary to work with superposition states. The most common approach consists in performing the Ramsey fringe experiment described in Section 2.1.3 for fine-tuning the drive pulse frequency.

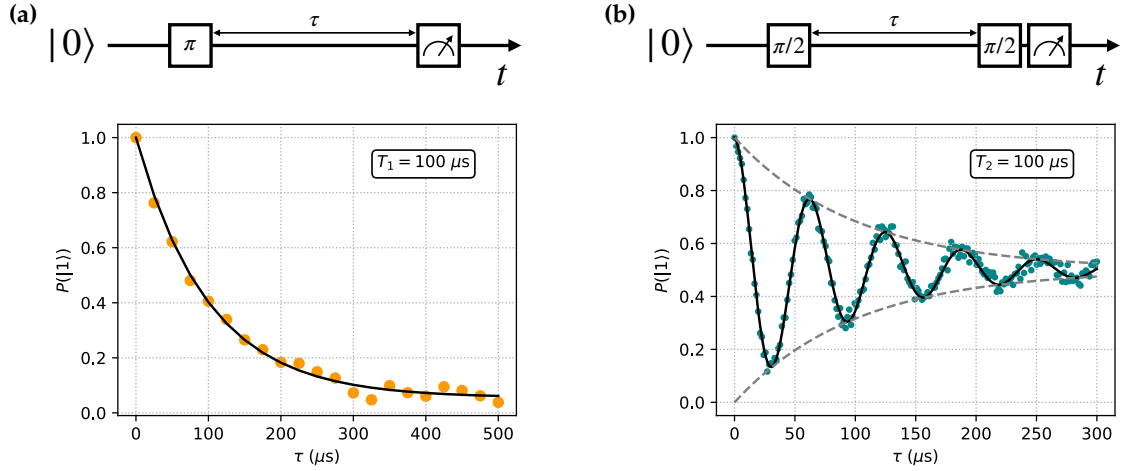


FIGURE 2.8: Simulations of (a) a T_1 and of (b) a T_2 measurement. The $\pi/2$ -pulse in (b) is slightly detuned from the qubit frequency in order to observe the oscillations in $P(|1\rangle)$. Oscillations are damped following an (approximately) exponential envelope.

In that section it was shown how, by sending a $\pi/2$ -pulse to a qubit in $|0\rangle$, waiting some time τ and then sending a second $\pi/2$ -pulse, the probability to measure the qubit in $|1\rangle$ after the second pulse oscillates as the delay τ increases:

$$P_{|1\rangle}(\tau) = \cos^2\left(\frac{\Delta\tau}{2}\right) = \frac{1}{2}[\cos(\Delta\tau) + 1] \quad (2.65)$$

Where Δ is the detuning of the pulses with respect to the qubit frequency. In the presence of dephasing, random fluctuations $\delta\omega$ in the qubit frequency alter the precession of the Bloch vector around the z axis, adding random phases to the cosine term. As a consequence, even by keeping τ and Δ fixed, $P_{|1\rangle}$ fluctuates. Measuring the qubit state several times, then, keeping τ fixed (similarly to what is done for T_1 measurements), results in a value of $P_{|1\rangle}$ that is actually the average of its fluctuations due to dephasing. The more τ increases, the more random phase fluctuations will accumulate between the two $\pi/2$ -pulses, causing a higher dispersion of the cosine values that determine $P_{|1\rangle}$. Eventually, the cosine term will average to 0, leading to $P_{|1\rangle}(\tau \gg T_2) \approx 1/2$, where no information about the phase can be inferred anymore.

Repeating this process, each time varying the value of τ , then, will result in the amplitude of the oscillations of $P_{|1\rangle}$ being damped as τ increases because of decoherence (see Fig. 2.8b). In first approximation, this damping goes as an exponential decay $\exp(-\tau/T_2)$. It can be proven, however, that in the presence of $1/f$ noise the envelope features also a Gaussian component, being $\exp(-\tau/T_1) \exp(-\tau^2/T_\varphi^2)$ [94, 95].

2.3.4 Decoherence in superconducting qubits

Most of what said until now about decoherence is valid not only for superconducting qubits, but also for other qubit technologies. Longitudinal and transverse relaxation are problems common to different fields, not only quantum computing, and even the definitions of T_1 , T_φ and T_2 were first introduced not for qubits but in Nuclear Magnetic Resonance (NMR) [96]. In this section the focus will move to the decoherence sources affecting specifically superconducting qubits. Most of them can be ascribed to two main mediators: *two-level systems* (TLSs) and *quasiparticles*.

Interactions with two-level systems are believed to be the main cause of decoherence in superconducting qubits [97]. TLSs have been first proposed more than fifty years ago to describe some thermal properties of amorphous solids [98], which may seem uncorrelated with the materials used in the realization of superconducting qubits, but are actually quite common in qubit chips: oxidation layers due to atmospheric exposure, Josephson junctions, dielectrics, all of them are typically amorphous, and, therefore, a source of TLSs. Despite having been introduced so long ago, the microscopic nature of two-level systems is still a matter of debate: present-day proposals include, among others, atoms (or group of atoms) tunneling between two equilibrium positions [99] (see Fig.2.9a), tunneling electrons [100, 101] and dangling bonds [102]. Nonetheless, all these proposals share fundamentally the same mathematical description, which is commonly referred to as the *Standard Tunneling Model* [103].

In the Standard Tunneling Model the TLS is described by an asymmetric double-well potential (see Fig.2.9b): in the scenario of a TLS being an atom in this potential, the first two energy levels of the system corresponds roughly to the atom being localized, respectively, in the left and in the right well³. This atom, though, can tunnel through the potential barrier and jump from one level to the other. In doing so, it has to exchange energy with the surrounding environment, emitting or absorbing energy quanta.

If the energy separation between the two TLS levels is the same of the energy separation between $|0\rangle$ and $|1\rangle$ in the qubit, the TLS can coherently exchange energy with it, resulting in longitudinal relaxation of the qubit. Nonetheless, even when the level separation of TLSs is very different from the one of the qubit decoherence may occur: TLSs in the oxide forming the Josephson junction can alter the potential barrier between the two superconductors and, hence, the critical current I_c of the junction [104, 105]. As can be seen in Eq.1.20, fluctuations in the critical current result in fluctuations in

³The wavefunctions of the corresponding Hamiltonian are actually hybridized, being a linear combination of the lowest eigenstates of the two wells. How much these wavefunctions are localized is determined by the energy difference between the two minima and by the potential barrier separating them. In general, higher is the energy difference between the minima w.r.t. the potential barrier, the more the wavefunctions are localized.

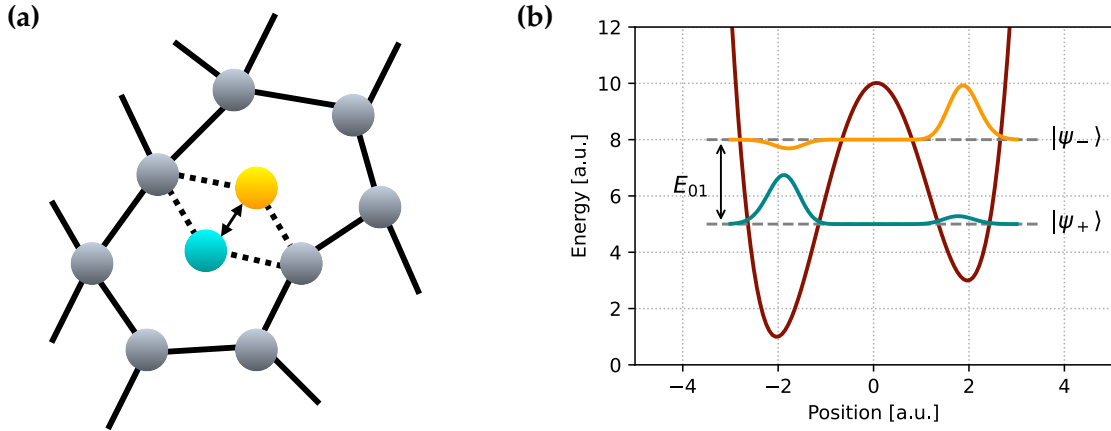


FIGURE 2.9: Example of a Two-Level System: (a) atom in an amorphous solid tunneling between two potential minima and (b) representation of the double-well potential with the wavefunctions of the two states.

the Josephson energy E_J and, consequently, in the qubit frequency ω_{01} (see eq.1.31), causing qubit dephasing [106]. TLS-TLS interactions are also possible, and are believed to cause fluctuations in the intensity of the noise affecting qubits, that can lead to jumps or drifts over time of the T_1 and T_2 values [107].

Due to the huge impact that TLSs have on qubit lifetimes, reducing their presence on chips and designing qubits that are insensitive to them is one of the main effort of researchers. As of today, mitigation strategies like adding capping layers to prevent surface oxidations [89] or using different superconductors [108, 109] have already been implemented, but the problem is far from being solved.

The second main source of decoherence, commonly named quasiparticles, is broken Cooper pairs. Quasiparticles are problematic because, when tunneling through the Josephson junction, can exchange energy with the qubit causing energy relaxation [110, 111]. Other than that, the low-frequency noise caused by the quasiparticle current tunneling through the junction affects the dephasing rate of the qubit, especially in flux-tunable transmons [112].

In an ideal experiment quasiparticles should not be a problem, because their density in a superconductor drops exponentially with the temperature [91]. At 20 mK, this value is supposed to be negligible. And yet, multiple experiments [113, 114] observed that the typical density of quasiparticles in a superconductor is much higher than that, with values ranging between 10 and 100 μm^{-3} , concluding that some other phenomena are at play.

Sources of quasiparticle poisoning are still under investigation, but some of them are already known. The first one to be identified were infrared photons [115, 116], which

presence is mostly due to the blackbody radiation of higher temperature stages in dilution refrigerators.

More recently, quasiparticles have been proposed as a mediator for vibration-induced qubit decays [117], which can be problematic especially in dry cryostats, where pulse tube coolers are known to be a source of vibrations.

Finally, in the next chapters it will be shown how also environmental radioactivity induces qubit decays by producing quasiparticles [118].

Several practices to mitigate the impact of quasiparticles on qubits have already been implemented, from shielding against external radiation [119] to the development of traps at the chip level [120, 121]. Due to variety of quasiparticle sources, it is likely that a combination of all these techniques will be needed in order to reduce significantly their impact on superconducting qubits.

Chapter 3

Qubits and Radioactivity

3.1 Radioactivity as a source of quasiparticles

Environmental radioactivity can be found virtually everywhere [122]. It can arise from the spontaneous decay of unstable nuclei in the laboratory environment or from contamination of the materials used to fabricate the experimental setup. It can also come from astrophysical sources, reaching Earth as cosmic rays. It can even be the residual product of nuclear accidents and weapon testing that occurred in the last century (^{137}Cs is a well-known example).

The emitted particles have, in the case of isotope decays, energies spanning from the keV scale up to a few MeV, while cosmic rays can reach energies several orders of magnitude higher due to multiple acceleration mechanisms in the Universe [123]. Typical energy deposits vary between sources: they range from the tens or hundreds of keV of high-energy photons, which could even cross a qubit chip without interacting at all, to the MeV of α particles that are, instead, completely absorbed. In any case, this amount is much higher than the binding energy of Cooper pairs, which is below 1 meV for the materials used in qubit fabrication [36]. It is straightforward to conclude, then, that an interaction with a qubit would raise the density of quasiparticles by several orders of magnitude with respect to its equilibrium value, with a disruptive effect on the coherence time of the device [124].

This was demonstrated in an experiment conducted at the Massachusetts Institute of Technology (MIT) in 2020 [118], in which a chip with two transmon qubits was placed in close proximity to a ^{64}Cu source. They then monitored the decay rate Γ_1 for both qubits over a few weeks.

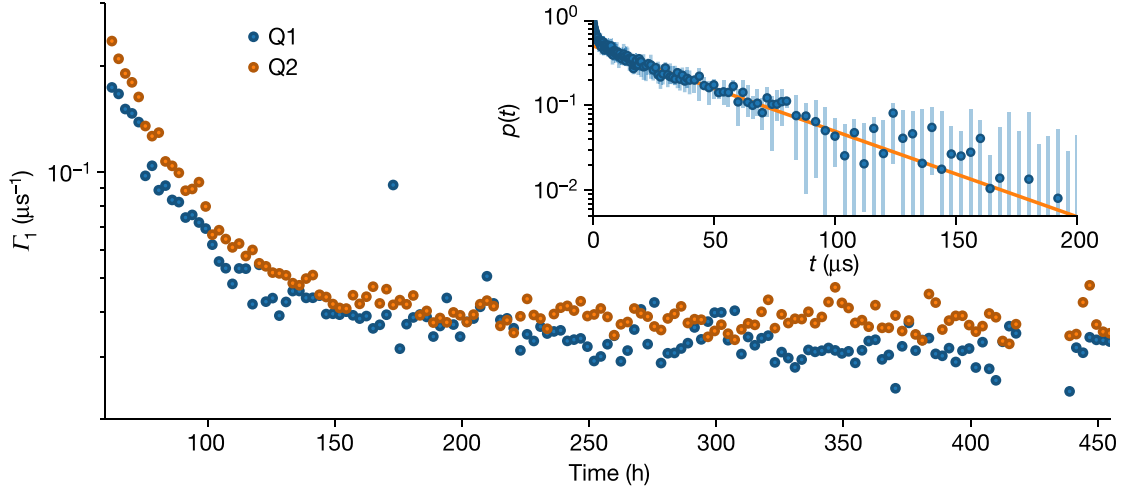


FIGURE 3.1: Decay rate of two qubits placed near a ^{64}Cu source as a function of the time elapsed since the start of the experiment. In the small plot, an example of a decay rate measurement, with $p(t)$ being the probability to measure the qubit in $|1\rangle$ a time t after having prepared it. Measurement strategy is the same described in section 2.3.3.

Plot from [118].

The ^{64}Cu source was chosen because it has a half-life of only 12.7 hours. Its activity therefore decreases over the measurement time, along with its contribution to the decay rates of the qubits. This is exactly what was observed (see Fig.3.1): decay rates were initially around $1/6\,\mu\text{s}^{-1}$ for both qubits, but then dropped to $1/35\,\mu\text{s}^{-1}$ within a week, by which time the ^{64}Cu source had almost completely decayed.

To corroborate this result, the research group repeated the measurement following the opposite approach: instead of exposing the chip to a radioactive source, they shielded it against environmental radioactivity. They built a movable 10-cm-thick lead shield around the cryostat and repeated the decay rate measurement several times with and without the shielding. As discussed in the next chapter, this type of shielding does not completely protect against environmental radioactivity (radioactive isotopes contaminating the setup had not been removed, and muons produced in cosmic-ray showers can easily penetrate through this layer of lead), but can still reduce the amount of radiation reaching the qubits, especially from environmental γ -rays that are produced by radioactive decays in the laboratory environment.

In this second experiment, however, the outcome was not as evident. The decay rate of one of the two qubits was on average lower when the chip was shielded compared to when the shield was off, but this difference was very small. The average $\delta\Gamma_1$ was indeed only $1/15.5\,\text{ms}^{-1}$, 0.26% of the total decay rate. This outcome has been interpreted as a proof that environmental radioactivity is not the most concerning source of decoherence in present-day qubits, whose dynamics are instead dominated by other decay

channels. However, researchers concluded that, as the performance of superconducting qubits will improve in the coming years, ionizing radiation will become increasingly relevant, requiring the development and implementation of mitigation strategies.

From their measurements, researchers also provided an estimation of how much ionizing radiation was affecting the density of quasiparticles in the qubits. This information is usually given in terms of the *normalized quasiparticle density*, which is defined as the ratio between the density of quasiparticles and the density of Cooper pairs in a superconductor:

$$x_{qp} = \frac{n_{qp}}{n_{cp}} \quad (3.1)$$

And can be estimated from the quasiparticle-induced decay rate by means of the equation:

$$\Gamma_{qp} = \sqrt{\frac{2\omega_q\Delta}{\pi^2\hbar}} x_{qp} \quad (3.2)$$

Where ω_q is the qubit frequency and Δ is the superconducting gap. From the collected data, the researchers estimated x_{qp} to be approximately $7 \cdot 10^{-9}$. Even if this value poses a limit on the coherence time of next-generation qubits, it would theoretically be possible to deal with the associated error rate by implementing quantum error correction. Error correction algorithms have indeed been successfully tested on qubits with decay rates much higher than the one that this quasiparticle density would induce [125].

3.2 The role of the substrate

By looking only at the MIT result, it would seem that ionizing radiation is a troublesome source of quasiparticles in superconducting qubits, but its effects could still be mitigated by quantum error correction. Unfortunately, other studies have shown that the interplay between radioactivity and qubits is more complicated than that, with other effects that make the implementation of error correction algorithms infeasible.

The first of these studies, Cardani et al [119], investigated the production of quasiparticles from environmental radioactivity in a chip with three superconducting resonators on it. The chip was characterized in three different locations: at the Karlsruhe Institute of Technology (KIT), Karlsruhe, Germany, at the University of Rome "La Sapienza", Rome, Italy, and at the Gran Sasso National Laboratory (LNGS), L'Aquila, Italy. In Karlsruhe and in Rome the chip was unprotected against environmental radioactivity,

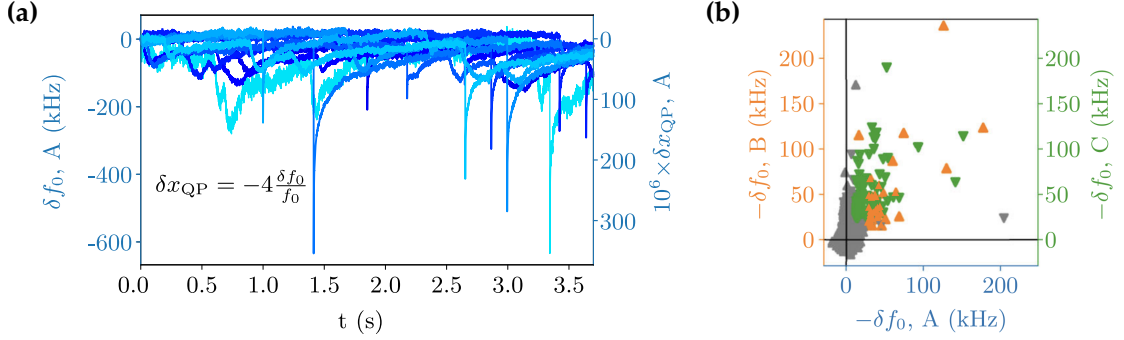


FIGURE 3.2: (a) Frequency shifts induced by quasiparticle bursts in one of the three resonators with the corresponding quasiparticle density. (b) On the y axis, frequency shift of events in resonators B and C coincident with an event in resonator A, with its corresponding frequency shift on the x axis. Plots from [119].

while in LNGS the experiment was run in a deep-underground facility designed for low-radioactivity experiments: the facility is placed below the Gran Sasso mountain, surrounded by 1.4 km of rock. The mountain above acts as a natural shield for cosmic rays, that could not be blocked in the MIT experiment, and is able to reduce the flux of cosmic muons by six orders of magnitude, down to $1 \text{ muon}/(\text{m}^2 \text{ h})$ [126]. This flux is practically negligible if considering typical superconducting chip prototypes, which have a $O(\text{cm}^2)$ surface. In order to shield the chip also from environmental γ -rays, the dilution refrigerator was surrounded with a 10-cm-thick lead shield.

The measurement strategy is analogous to the one used for *kinetic inductance detectors* (KIDs), which are also made of superconducting resonators [127]. There, a quasiparticle burst in the superconductor changes the kinetic inductance of the resonator, producing a shift in the resonant frequency δf_0 proportional to the fluctuation in the density of quasiparticles δx_{qp} :

$$\delta x_{QP} = -\frac{4}{\alpha} \frac{\delta f_0}{f_0} \quad (3.3)$$

Where α is the ratio of kinetic inductance over the total inductance. Measurement is then done by probing the resonator at f_0 and looking for shifts in its response induced by quasiparticle bursts (see Fig.3.2a).

From the last equation it is also possible to see that, if the ratio of kinetic inductance is small, a quasiparticle burst will produce a small shift in the resonant frequency, complicating the measurement. For this reason, the resonators used in the study were made of granular aluminum (grAl), for which $\alpha \approx 1$ [128].

By comparing the results obtained in the shielded configuration with what was achieved in the other two locations above-ground, it was confirmed that environmental radioactivity produces quasiparticle bursts, but two major complications for quantum error correction emerged.

The first one, as can be seen in Fig.3.2a, is that in a single burst the quasiparticle density can rise up to 10^{-4} before recovering, a value orders of magnitude higher than the $\approx 10^{-8}$ obtained by the MIT experiment by looking at longer timescales. Even though the event reconstruction showed that this density is transient, with Cooper pairs being formed again within seconds from the particle interaction, it is still critical for quantum computation. No present-day algorithm for quantum correction would be able, in fact, to tolerate the error rate induced by such a density on a qubit, resulting in the failure of every computation going on when the impact happens.

The second complication could be even worse for error correction. It was found, in fact, that frequency shifts of the three resonators were correlated in time (see Fig.3.2b), meaning that all of them were affected by the same particle interaction. This happens because the particle interaction does not need to be located in the superconductor to affect the resonator (or the qubit), which is actually unlikely due to its very small dimensions. The interaction can actually take place in the substrate of the chip, which is much larger and thicker than the superconductor, producing ballistic phonons. Phonons then propagate throughout the chip, reaching all the resonators and breaking Cooper pairs. Phonon-mediated quasiparticle poisoning is a mechanism that has actually been observed in superconducting circuits [129] (it was also suggested, but not experimentally confirmed, as a decoherence channel in the MIT study), and has been extensively studied with kinetic inductance detectors [130].

This result is detrimental when working with superconducting qubits, because, as discussed in the previous chapter, error correction works only if errors are uncorrelated. If environmental radioactivity is capable of making multiple qubits fail simultaneously, then, it poses a great obstacle in the realization of a fault-tolerant quantum computer.

3.3 The problem of correlated errors

3.3.1 The first observation

The confirmation that radioactivity can induce correlated errors in multi-qubit chips came from another research article, Wilen et al. [131], that was published almost simultaneously with the previous one. Radiation-induced correlated errors were observed by

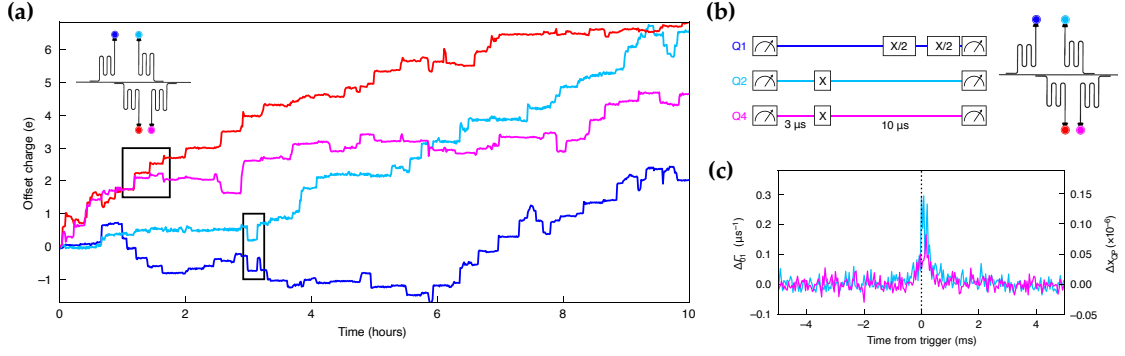


FIGURE 3.3: (a) Offset charge as a function of time for the four qubits; in the black boxes, two examples of correlated jumps. (b) Measurement scheme for the correlated charge jump and quasiparticle burst measurement. (c) A quasiparticle burst in two qubits correlated with a charge jump in the third one. Plots from [131].

monitoring the charge offset of four transmon qubits over time.

As described in Section 1.2.2, transmon qubits are designed to be charge-insensitive, so in a typical transmon radiation-induced fluctuations of the charge offset would be difficult to measure. The qubits used in this experiment, then, were engineered in order to have a E_J/E_C ratio smaller than what is usually chosen for transmon qubits, making them weakly charge-sensitive.

In a charge-sensitive qubit the presence of an offset charge alters its energy levels, eventually affecting the transition frequency ω_{01} . Calling n_g the offset charge (expressed as number of Cooper pairs) and $\bar{\omega}_{01}$ the charge-averaged transition frequency, the transition frequency of the qubit goes as [132]:

$$\omega_{01} = \bar{\omega}_{01} + \Delta\omega_{01} \cos(2\pi n_g) \quad (3.4)$$

With $\Delta\omega_{01}$ dropping exponentially with $\sqrt{E_J/E_C}$ [30]. For the qubits used in the experiment $E_J/E_C = 24$, resulting in $\Delta\omega_{01}$ being $O(\text{MHz})$.

As shown in section 2.1.3, the best way to measure the transition frequency of a qubit is by performing a Ramsey fringe experiment. Researchers then repeated over time Ramsey measurements of all four qubits simultaneously. From ω_{01} they derived the offset charges, observing that there were events in which multiple qubits were experiencing similar abrupt fluctuations at the same time (see Fig.3.3a). These fluctuations, corresponding to jumps in the values of n_g , were more frequent when qubits were closer in space, so when they were on the same side of the feedline used to send readout pulses.

To prove that these charge jumps were induced by radioactivity, researchers looked for coincident quasiparticle bursts similar to the ones observed in the experiment with the

grAl resonators. In order to do so, with one of the qubits they kept doing Ramsey measurements looking for jumps in its charge offset. At the same time, though, they probed the decay time of other two qubits by exciting them and then measuring their state after a few μs (the idea being that higher is Γ_{01} , lower will be the probability to measure them in $|1\rangle$ after the time delay). Fig.3.3b shows the measurement scheme of this experiment.

In doing so, they observed that some of the charge jumps in the first qubit were associated with quasiparticle bursts in the other two qubits, responsible for drops in their decay rate (see Fig.3.3c). This result points to the same mechanism observed in the experiment with resonator, with radioactivity producing phonons in the chip substrate that then propagate, reaching all the qubits and breaking Cooper pairs.

Researchers then performed a Monte Carlo simulation to estimate the contribution of radioactivity to these quasiparticle bursts, concluding that environmental γ -rays and cosmic muons produce a rate of events which is approximately 10 mHz, around half of the one measured. They also hypothesized that the remaining events could be ascribed to other sources like contaminations in the setup materials.

3.3.2 A real case study: correlated errors in the Google Sycamore

The chip used by Wilen et al. was engineered in order to be more sensitive to charge noise than a typical transmon, which could lead to the conclusion that that in a standard chip qubits are protected by design against correlated errors. Unfortunately, this is not true, as correlated errors were observed also in other devices. It is the case of the Google Sycamore processor, which was already used to prove that quantum computers are capable of performing calculations orders of magnitude faster than classical computers [17].

Following the result obtained on the charge-sensitive qubit by the other team, Google researchers developed an improved detection strategy for correlated quasiparticle bursts, described in McEwen et al. [133]. The idea was still to prepare qubits in $|1\rangle$ and then measure its state after a few μs , but this time there was no offset charge measurement for triggering, and the state preparation was improved in order to reduce the times the qubit was accidentally prepared in the wrong state.

This is because, by repeatedly sending π -pulses to qubits in order to keep them in $|1\rangle$, if they have not decayed before the new pulse, the new pulse will set them in $|0\rangle$ instead. In order to be sure that the final state is $|1\rangle$, then, consecutive cycles of state preparation and measurement must be separated by a time delay larger than T_1 . This means that the

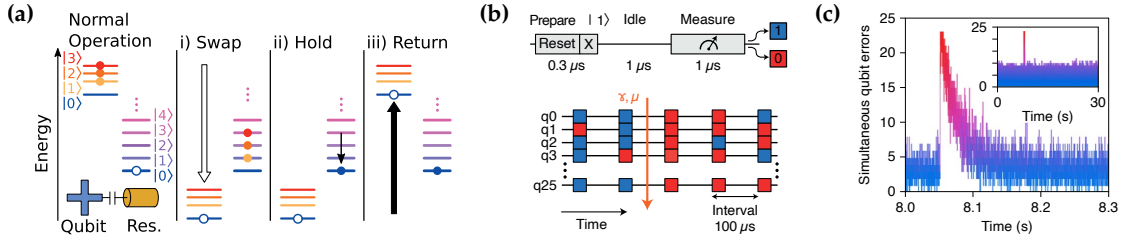


FIGURE 3.4: (a) The Google reset protocol, (b) the measurement strategy used to find correlated errors and (c) plot of an event with the number of qubits simultaneously decaying versus time. Image (a) is from [134], while (b) and (c) are from [133].

sampling rate cannot be too fast, which undermines the sensitivity of the experiment. Vice versa, by choosing qubits with small T_1 , so to be sure of their decay after a short time, the rate of errors caused by other decay mechanism will rise accordingly, making radiation events more difficult to identify.

Google researchers overcame this limitation by *resetting* [134] the qubits in $|0\rangle$ before sending the new π -pulse. The reset strategy implemented relies on the fact that all qubits are flux-tunable: in the first step of the reset their transition frequency gets reduced, going from higher than the ones of the readout resonators to lower than them. When the frequency of a qubit reaches the one of its resonator, the two elements can exchange energy, so excitations are swapped between them and the qubit decays to $|0\rangle$. After that, the qubit frequency is kept lower than the resonator one, allowing for the resonator to decay, and then is brought back to the original value (see Fig.3.4a). Since the decay time of the resonator is much lower than the qubit T_1 , this method allows to dissipate excitations much faster than what would be possible by just waiting for the qubit to spontaneously decay. In the Google chip, the entire process lasts less than $1 \mu\text{s}$.

Once the qubits were prepared in $|1\rangle$, their state was measured again after $1 \mu\text{s}$. That delay was chosen to be much smaller than T_1 , so in normal condition the outcome of the subsequent measurement was still $|1\rangle$. In the presence of quasiparticle poisoning, though, the decay time drops and the probability to measure $|0\rangle$ after the delay becomes significant.

The outcome of the experiment is shown in Fig.3.4c: in typical conditions, out of the 25 qubits measured, only a small fraction of them, between 0 and 5, would be found in $|0\rangle$. There were moments, though, in which almost all of the qubits decayed simultaneously, resulting in a general failure of the chip. The time separation between these events followed a Poisson distribution, with an average rate of one event every 10 seconds, similar to what was observed in the other experiment.

By looking at the positions of the decaying qubits, researchers also observed that these events started with a localized group of qubits decaying and then propagated across the

entire chip. This behavior is compatible with a phonon-induced burst of quasiparticles starting in the qubits closest to the interaction point, and then propagating to the other qubits as phonons spread over the chip, in agreement with the previous experiments.

3.3.3 Recent confirmations

After the Google result, new confirmations that ionizing radiation is responsible for at least a fraction of these phonon-induced quasiparticle bursts arrived very recently from three different studies. These studies were all published in the final part of my PhD project, between May and November 2025, but it is worth mentioning them since they reached conclusions similar to the ones of my thesis.

In the first one, researchers repeated the experiment described in Section 3.3.1, where correlated jumps of the offset charge were observed. This time, though, the chip used in the first experiment was measured at Fermilab, US, with different levels of shielding against environmental radioactivity [135]. To abate the flux of cosmic muons, the setup was placed in an underground facility, approximately 100 m below Earth's surface. Then, measurements were performed with and without a lead shield blocking environmental γ -rays. Both the rates of events obtained, with and without the lead shielding, are two orders of magnitude lower than what was measured above ground in [131], confirming that cosmic rays are among the causes of these jumps. However, by comparing the two configurations it was observed that the lead shielding reduced the rate of correlated jumps only by a bit less than a factor 3, from 0.51 mHz down to 0.19 mHz. This reduction is much smaller than what would be expected if the events observed underground were caused mainly by γ -rays, pointing out to other sources of quasiparticles not yet taken into account.

The other two studies used an auxiliary particle detector to identify events generated by muon interactions. In the first one [136], two scintillators, i.e. materials that emit light when there is a particle interaction inside them, were placed right below the qubit chip (see Fig. 3.5a). To observe the light emitted in an interaction, the scintillators were coupled to silicon photo-multipliers. If a cosmic muon crosses the superconducting chip with Zenith angle small enough, it will cross also the two detectors below it, producing a coincident signal in the three devices. This strategy is not effective when trying to detect γ particles, though, because, not having electric charge, they do not necessary interact in all the devices.

The second study [137] followed a similar principle, using again scintillating detectors. This time, though, the detectors were placed farther from the chip, below the dilution refrigerator used to cool it down (see Fig. 3.5b).

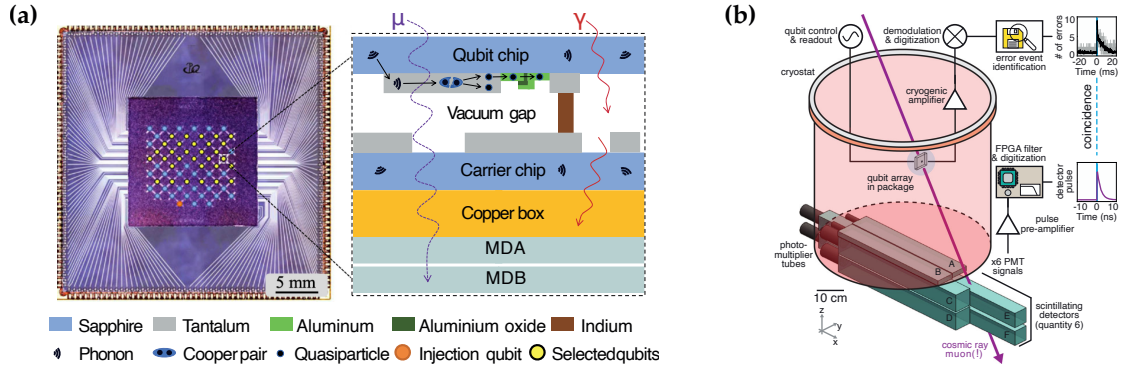


FIGURE 3.5: Schemes of the setups for the detection of muon-induced quasiparticle bursts (a) in [136] and (b) in [137], taken from the respective articles.

In the first study, researchers followed two different measurement strategies. First, they performed repeated Ramsey measurement to detect jumps in the offset charge, in a way similar to what was done by Wilen et al [131]. For qubits with low E_J/E_C ratio, this type of measurement is the most sensitive one to quasiparticle bursts, because even the tunneling of a single quasiparticle through the junction results in a measurable variation of the qubit frequency. Then, they also looked at sudden increases in the qubit decay rate by using a protocol analogous to the one described in McEwen et al [133].

In both cases, they were able to time correlate muons detected with the two scintillators with quasiparticle bursts in the chip, confirming that cosmic rays are responsible for at least a fraction of the bursts observed.

The second study, instead, focused only on the decay rate of the qubits, using the same measurement strategy of McEwen et al [133]. This time, no qubit reset at the beginning of the cycle was implemented. Nonetheless, researchers managed to successfully identify muon-induced quasiparticle bursts, concluding that cosmic rays account for approximately 17% of the bursts observed.

They also observed that muons are responsible mostly of events involving a higher number of qubits. This was not the most common type of event, where instead only 4 or 5 qubits were involved. They hypothesized that events involving a smaller number of qubits were caused by γ particles, which on average deposit a smaller amount of energy in the substrate.

Both these works confirmed the contribution of cosmic muons to quasiparticle bursts in superconducting devices. Building on that, our research group proposed the realization of a muon tagging system for the identification of muons interacting in the qubit chip. To maximize the efficiency of this system, we decided to rely on cryogenic detectors that could be placed close to the chip, reducing the rate of false triggers. The system consists of two kinetic inductance detectors, placed respectively above and below the quantum

chip. This system was tested for the first time in February 2025, using a third KID as a proxy for a quantum device, achieving a muon tagging efficiency of 90%. Results have been summarized in Mariani et al [138], currently under peer review for publication. The entire article is reported in Appendix A.

3.4 Impact on Two-Level Systems

All the results discussed until now considered only scenarios in which ionizing radiation directly affects superconducting qubits. However, phonons produced in a particle interaction propagate all over the chip, not only towards the qubits. It is possible, then, that ionizing radiation indirectly changes the qubit behavior by modifying the surrounding environment, for example by interacting with two-level systems. This phenomenon was first observed by IBM researchers, who named it *TLS scrambling* [139].

In their study, they monitored 17 transmon qubits on the same chip, looking both for correlated jumps in the charge offset and for transient increases in their decay times. The transmon qubits used in the experiment were not engineered for a low E_J/E_C ratio, so the frequency dispersion with the charge was much smaller than what had in Wilen et al. [131] or in Li et al. [136]. To overcome this limitation, Ramsey spectroscopy was performed between $|1\rangle$ and $|2\rangle$, which are more sensitive to charge fluctuations.

Alternated with Ramsey spectroscopy there was a T_1 measurement following the same excite-wait-measure protocol used in the Google experiment. The overall measurement then consisted in:

- Prepare the qubit in $|1\rangle$;
- Perform Ramsey spectroscopy between $|1\rangle$ and $|2\rangle$ and measure the outcome;
- Prepare again the qubit in $|1\rangle$;
- Measure again after $40\ \mu\text{s}$.

Taking approximately $44\ \mu\text{s}$ to complete a full cycle.

Researchers observed simultaneous charge jumps of qubits separated also by a few millimeters, that allowed them to exclude local sources of charge noise in favor of environmental radioactivity. The rate of these events was tens of mHz, compatible with what observed also by other experiments.

Concerning the decay time, also in this experiment sudden increases, correlated with a charge jump, were observed. This time, though, for all the qubits the decay time was

recovering faster than what observed on other chips, probably due to different choices of materials and geometry of the qubits.

By looking at the T_1 values, researchers observed also that, after some events, T_1 's were stabilizing to values different from their averages before the impact, suggesting that something else was at play.

Typically, these fluctuations of a qubit's lifetime are associated with modifications of the two-level system environment coupled to it, a phenomenon that was already mentioned in section 2.3.4. To prove that this is true also in this case, though, it is necessary to do a *TLS spectroscopy*. It is a technique relying on flux-tunable qubits, whose transition frequencies are swept over their full range. At each frequency step, a T_1 measurement is done [107]. The idea behind this approach is that, when a qubit is at resonance with a TLS, its T_1 drops. It is then possible to find the frequencies of TLSs coupled to the qubit by looking at its T_1 values at different frequencies.

The main obstacle in doing a TLS spectroscopy also for this chip was that the transmon qubits used were not flux-tunable. Researchers overcame this issue by inducing an *AC Stark shift*: it is a shift in the energy levels of the qubit caused by an off-resonance driving. This alters the transition frequency, allowing for small sweeps even if the qubit studied is not flux-tunable. It is a technique that has been used already in the past to investigate TLSs in proximity to a qubit transition frequency [140].

The shift induced with this technique on the qubits used for the IBM study allowed to span only over a few tens of MHz, but that was already enough to confirm that particle interactions were indeed affecting TLSs. Researchers alternated Ramsey measurements, necessary to detect particle interactions, and TLS spectroscopy done with the Stark effect, looking for changes in the TLS spectrum induced by the same particle interaction, the TLS scrambling named at the beginning of this section.

Fig. 3.6 shows three examples of TLS scrambling: on the top panels the correlated charge jumps of different qubits are visible, pointing out to a particle interaction in the chip. Simultaneously with that, the TLS spectrum changes abruptly, as can be seen in the middle panels: on the y axis is the shift in the qubit frequency induced with the Stark effect, while on the x axis is time; the colors represent the probability to measure the qubit in $|1\rangle$ after the $40\ \mu\text{s}$ delay, which depends on T_1 . Lowest probabilities are associated with the presence of TLSs at that frequency. It can be seen that the positions of the most evident TLSs change after the impact.

To quantitatively measure TLS scrambling, researchers introduced Pearson's r in their analysis. It represents the linear correlation between two sets of ordered data: when $r = 1$ the correlation is maximum, meaning that the two datasets are identical; the more

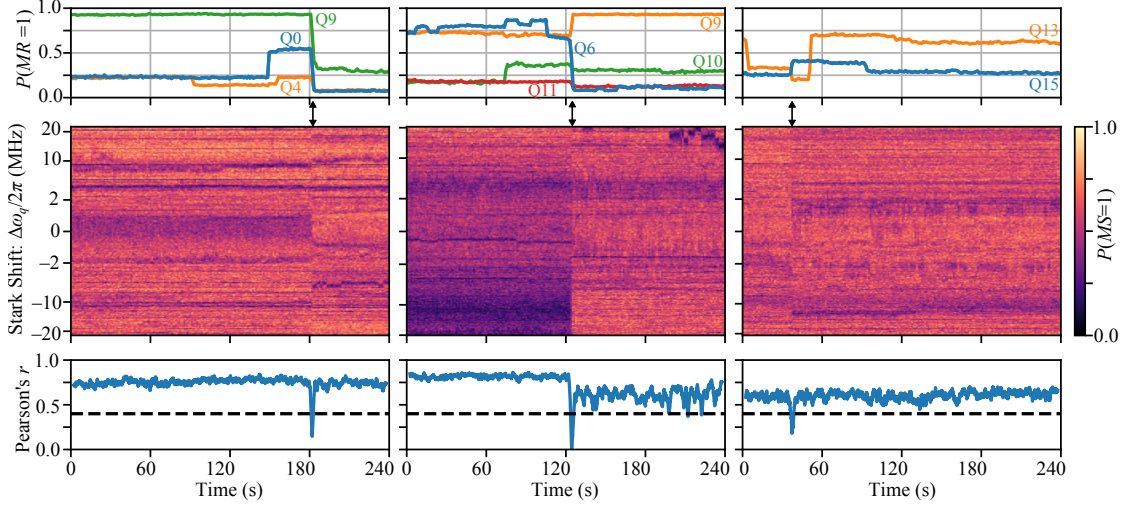


FIGURE 3.6: Example of three events where the particle interaction, detected by the correlated charge jump of multiple qubits in the top panels, is followed by a change in the TLS spectroscopy (middle panels). In the bottom panels, Pearson's r of the TLS spectroscopy data, with 0.4 used as a threshold value to identify a scrambling event.

Image from [139]

the datasets differ, the more r gets closer to 0. In the case of this study, Pearson's r was used to compare the TLS spectrum right before and right after particle impacts. High values of r meant that the TLS spectrum was not affected, while low values indicated that TLS frequencies changed after the event. Pearson's r for the three example events in Fig. 3.6 is shown in the bottom panels. The dashed line corresponds to a value of 0.4, used as a threshold to identify scrambling.

The discovery of TLS scrambling was accompanied with new open questions, particularly concerning the rate of the events with associated scrambling. This was found to be much lower than the total rate of particle interactions, with less than 0.1% of the charge jump events followed by a TLS scrambling observed by multiple qubits. Researcher hypothesized that this can be a consequence of the dipole moment of two-level systems, which is, in most models, orders of magnitude smaller than the one of a qubit. As a consequence, two-level systems are less sensitive to the charge redistribution following a particle interaction, and the scrambling could involve only a small, localized fraction of the TLSs. This hypothesis, however, still lacks an experimental proof.

3.5 Summary and outlook

Studies performed in the past years unequivocally proved that superconducting qubits are sensitive to environmental radioactivity: impinging particles affect the coherence time of qubits [118] and the overall performance of superconducting devices [119], are a

source of correlated errors in multi-qubit chips [131, 133, 135–137] and can even affect qubits indirectly through TLS scrambling [139].

The solutions proposed for this problem are many. A lot of them aim at protecting qubits from phonons produced in particle interactions. A non-exhaustive list of proposals of this kind includes adding *phonon traps* to quantum processors, depositing qubits on modular dies isolated from each other and engineering the superconducting gap to make them insensitive to ballistic phonons.

Phonon traps are one of the earliest proposals for phonon mitigation, being already investigated for space telescopes based on kinetic inductance detectors [141]. They consist in thin superconducting layers with low critical temperature, designed to absorb phonons and down-convert their energy to values below the superconducting gap of qubits [142, 143].

The modular fabrication of quantum processors has been instead proposed by Rigetti Computing [144]. Their demonstrator consisted of four eight-qubit integrated circuits, flip-chip bonded to a larger carrier chip. This type of fabrication improves the isolation of different regions of the processor compared to depositing all qubits on the same substrate. In this way, radiation-induced phonons could be more easily confined into a portion of the chip, reducing the risk for a general failure of all the qubits.

Finally, gap engineering has been recently proposed by Google as a method for reducing qubit sensitivity to phonons [145]. The idea behind it can be seen as the opposite of phonon traps: while phonon traps down-convert phonon energy to make it lower than the superconducting gap, gap engineering designs qubits such that their superconducting gap is higher than the energy of the phonons produced by ionizing radiation.

All these proposals showed promising results in protecting qubits from environmental radioactivity, but none of them fully resolves the problem.

A complementary strategy, that could work in synergy with the proposals just described, consists in shifting the focus to the radioactivity sources most affecting superconducting qubits, aiming at mitigating their impact. This is the approach followed by our research group in this work. The first step for the implementation of this strategy consisted in understanding which radiation sources are of most concern in a typical qubit experiment. This was done by developing a Monte Carlo simulation, and will be the topic of the next chapter.

Chapter 4

The radiation background of superconducting qubit devices

As seen in the last chapter, superconducting quantum bits are affected by several sources of ionizing radiation, which degrade qubit coherence and induce correlated errors that cannot be corrected with present-day algorithms. Experiments conducted before my PhD demonstrated this impact, though it was broadly attributed to cosmic rays, with other sources being overlooked. Determining which of them are most impactful for superconducting qubits, though, is not straightforward. Each source produces ionizing radiation with different intensities, but their contributions to qubit errors also depend on how far particles are produced from the chip, the material between the sources and the qubits, and the type of radiation produced. As an example, α particles produced by the decay of ^{222}Rn , a radioactive isotope commonly found in air, do not even reach a chip inside a dilution refrigerator, since the metal vessels between temperature stages can easily block them. In contrast, the same vessels do not block γ particles produced by the daughter isotopes of ^{222}Rn , which are therefore much more concerning.

Measuring the intensity of radiation sources is thus only the first step to estimate their impact on qubit experiments. It is also necessary to estimate how many of the particles produced actually reach the chip and interact with it. These last two pieces of information, the path of the particles and their interactions in the chip, are intrinsically probabilistic, since particle dynamics is quantum mechanical. The easiest way to account for these phenomena is to employ a Monte Carlo simulation, which is particularly useful when working with complex geometries, such as the experimental setup of a qubit experiment.

The first part of my PhD was thus devoted to developing such simulation and to estimating the contributions of the main radiation sources to qubit coherence. The simulation

was built with Geant4 [146–148], a well-validated toolkit developed by CERN in collaboration with other research institutions, and currently the most widely used simulation toolkit in particle physics. The results of the simulation have been reported in an article titled *Disentangling the sources of ionizing radiation in superconducting qubits* [149], published on the European Physical Journal C in January 2023.

4.1 Sources of ionizing radiation

Ionizing radiation has historically been one of the major problems of particle physics experiments looking for rare events, like interactions of neutrinos, dark matter detection, or hypothetical nuclear decays [122]. Particles from environmental radioactivity that interact in the detectors produce a background of false positive events in the experiment data that makes it difficult to identify the signature of the physical process under study. If the rate of these events is too high compared with the rate of the events of interest, no meaningful conclusion can be drawn from the data. Some of the background events can be recognized as such and discarded: this is usually done by *vetoing*, i.e. surrounding the experiment with auxiliary detectors to identify incoming particles from environmental radioactivity, or by looking at the event *topology*, i.e. the response of the detector (or the detectors, if more than one are involved) to different particles, different types of interaction or different energies and momenta involved. Even after these procedures, however, there is an irreducible background of events mimicking the signature of the process under study that cannot be eliminated, undermining the sensitivity of the experiment.

Since radioactivity can be so detrimental for these experiments, researchers have produced an extensive scientific literature about the different radiation sources, the way different particles interact with matter and the best mitigation strategy against each source.

In low-radioactivity experiments it is common to distinguish between *far* and *close* sources, depending on the different mitigation strategies that have to be implemented against them. Far sources are those sources from which the experiment can be shielded, such as particles produced by nuclear decays in the laboratory or coming from cosmic rays. On the contrary, close sources are located in the proximity of the experiment or are part of the experiment itself, like unstable isotopes contaminating the setup materials. Because of their location, shielding is ineffective against them, and different strategies, like the choice of purer materials, have to be considered.

Among far sources, environmental γ -rays are the most prominent. They are emitted after a radioactive decay, when the daughter nucleus is produced in one of its excited states and subsequently relaxes to the fundamental state [150]. γ -rays are therefore always accompanied by α or β particles. These other products, though, have a much shorter range, being absorbed by small amounts of solid materials (α particles can be blocked with a sheet of paper, while electrons and positrons penetrate at most a few centimeters in metals). Because of that, they play a role only when produced in the vicinity of the device, and their contribution as far sources can be safely neglected.

The other major far source is cosmic radiation, commonly referred to as cosmic rays. Particles coming from outer space are mostly high-energy protons, but when crossing Earth's atmosphere they undergo a cascade of nuclear reactions, producing a variety of secondary particles like neutrons, electrons, muons, neutrinos and other protons. Secondary particles also interact with the atoms of the atmosphere, and the majority of them get absorbed before reaching the ground. As a result, at sea level more than 60% of incoming particles are muons, by far the most penetrating ones. Neutrons account for a little over 20%, while the remaining particles are mostly electrons and, less likely, protons [122].

Neutrons from cosmic rays represent most of the total neutron budget, but a residual neutron flux can be observed also deep underground, where cosmic rays do not arrive. These neutrons are produced by two nuclear reactions: the spontaneous fission of heavy nuclides and (α, n) reactions, where a neutron is emitted after the collision of an α particle with a nucleus [151].

As for close sources, most of the impurities can be ascribed to ^{40}K or to the progenies of ^{232}Th , ^{235}U and ^{238}U . These isotopes were produced in astrophysical processes before the formation of Earth: the reason why they can still be found in nature is that they have extremely long half-lives, from 700 million years of ^{235}U up to 14 billion years of ^{232}Th , so they have not completely decayed yet.

^{232}Th , ^{235}U and ^{238}U are the long-lived parents of three distinct decay chains. None of them decays directly into a stable nucleus; rather, each decays into another radioactive isotope, which will also decay and so on, until a stable isotope is reached. In this way tens of different radioactive isotopes are produced.

When the number of radioactive nuclei is constant, that is, when the decay rate of each isotope equals the decay rate of its parent, the chain is in *secular equilibrium*. Under these conditions, the overall rate of decays from all isotopes is constant, and, as a consequence, also the rate of impacts of their decay products in the qubit chip.

There are two other types of radioactive isotopes that have to be taken into account: anthropogenic ones (especially ^{137}Cs), produced by nuclear tests and accidents over the last century, and isotopes produced through *cosmogenic activation*. It consists of spallation processes caused by secondary cosmic particles impacting materials used in the fabrication of the experiment, especially when transported by air. The nature of these isotopes, along with their abundance, depends on the material being exposed to cosmic radiation. In qubit experiments, the most widely used material in proximity of the chip is copper, due to its very good thermal conductivity at low temperatures (necessary to properly cool down the device) [25]. For copper the most common cosmogenic isotopes are the Cobalt ones, especially ^{60}Co , ^{58}Co and ^{57}Co , with minor contributions also from ^{54}Mn and ^{59}Fe [152].

4.2 Measuring the activity of radioactive sources

As mentioned in the introduction of this chapter, the first step to understand which sources are more concerning for superconducting qubits is to quantify their intensity. For far sources, this is expressed in terms of flux of particles per unit surface and unit time. For close sources, instead, the relevant quantity is their *activity*, i.e. the number of radioactive decays per unit time. Since a larger mass of material contains a higher number of contaminating isotopes, the activity is measured in terms of decays per unit time per unit mass of material.

Measurements were done as a first step for the Round Robin experiment, done in the framework of the Superconducting Quantum Materials and Systems (SQMS) Center and aiming at characterizing the same multiqubit chip in different facilities. The chip design and fabrication is done by Rigetti Computing, while locations involved, and so chosen for radiation measurements, are Rigetti Computing headquarters, in California, the National Institute of Standards and Technology at Boulder (NIST), in Colorado, Fermilab and the Northwestern University, in Illinois, and the Gran Sasso National Laboratory (INFN-LNGS), in Italy.

4.2.1 Far sources

In measuring far sources, we focused on the most significant ones, so cosmic muons, neutrons and environmental γ -rays. The magnitude of their fluxes depends on the location and the structure of the laboratory, so for all of them different locations were considered, both above ground and deep-underground.

At sea level, the **muon** flux is about $1 \mu/(\text{cm}^2 \text{ min})$, with an average muon energy of $\approx 4 \text{ GeV}$ [153]. For muons with this energy, the angular distribution goes as $\cos^2 \theta$, where θ is the zenith angle. This is due to the fact that muons incoming from higher angles have to make a longer path before reaching ground, and ends up being more attenuated by Earth's atmosphere. For the same reason, the flux of muons gets higher as the altitude increases.

The biggest variation in the muon flux, though, is achieved by going deep underground. There, surrounding rock attenuates the muon flux much more than what the atmosphere does above ground, with the possibility of damping it by several orders of magnitude by going deep enough. The magnitude of this attenuation depends on the chemical composition of the rock, so, when comparing different underground locations, instead of the depth it is preferred to use the *meters water equivalent* (m.w.e.), i.e. how deep underwater should the experiment be placed in order to have the same attenuation. At the Gran Sasso Underground Laboratory, which is below 1.4 km of rock, the muon flux is attenuated by a factor 10^6 , corresponding to 3800 m.w.e. [126].

Following the same idea, also the location of the laboratory within its building, so whether it is in the basement or on one of the higher floors, can affect the muon flux, since the building itself can be seen as a shield. This variation, though, is minimal and often negligible.

The building plays a more significant role when considering **neutrons**. This is because neutrons lose energy mainly via inelastic scattering, while muons interact mostly via electromagnetic interactions. As a consequence, neutrons do not require hundreds of meters of material in order to be blocked, and their flux can vary up to a factor 2 between different locations.

Measurements of the neutron flux and spectrum were performed by using a DIAMON spectrometer [154]: it consists of a polyhedron of high-density polyethylene, which is used to absorb incoming neutrons, with semiconductor-based detectors to identify interactions. The choice of polyethylene as target materials is due to the low masses of its atoms, which grant a higher cross section for inelastic scattering with neutrons.

Neutron fluxes were measured in different laboratories, both above-ground and deep-underground. Concerning above-ground laboratories, fluxes ranged from $0.010 \text{ n}/(\text{cm}^2 \text{ s})$ up to $0.018 \text{ n}/(\text{cm}^2 \text{ s})$, so an average value of $(0.014 \pm 0.004) \text{ n}/(\text{cm}^2 \text{ s})$ was chosen for all locations. Fig.4.1a shows the energy spectrum measured above-ground at INFN-LNGS. Similar spectra were obtained also in other above-ground locations.

The situation is very different underground at INFN-LNGS, where the contribution of cosmic neutrons is orders of magnitude lower compared to the one of neutrons from

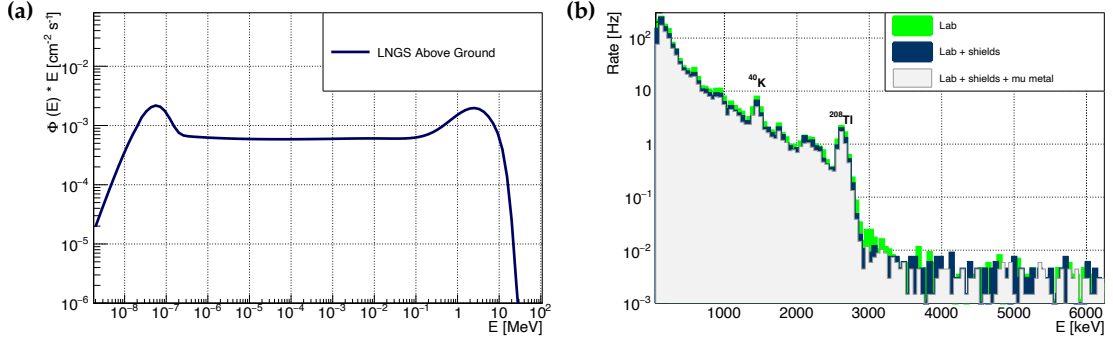


FIGURE 4.1: (a) Energy spectrum of neutrons measured above-ground at the Gran Sasso national laboratory and (b) energy spectrum of environmental γ -rays out of the dilution refrigerator (green), inside it (blue) and when the dilution refrigerator is surrounded by a mu-metal magnetic shield.

spontaneous fissions and (α, n) reactions. There, the total neutron flux was already measured by other experiments [155], finding a value that is four orders of magnitude lower than above-ground, around $10^{-6} \text{ n}/(\text{cm}^2\text{s})$.

Finally, environmental γ -rays are the only type of radiation whose flux does not vary by orders of magnitude between above-ground and underground facilities. This is because environmental γ -rays are produced by the decay of radioactive isotopes in the laboratory environment, instead of being the result of primary cosmic particles interacting with the atmosphere (see Section 4.1). Their energy spectrum is characterized by several peaks in correspondence to the decay energies of the isotopes present in the environment. The spectrum shows a rate drop at 2.615 MeV, corresponding to one of the decay peaks of ^{208}Tl ; above this energy, considered the end-point of environmental radioactivity, only rare γ -rays are observed.

The flux and spectrum of γ -rays was measured with a 3-inch portable NaI spectrometer. NaI is a scintillating material, so γ -rays are detected by the scintillation light produced when they interact with the crystal. Measurements above-ground at INFN-LNGS resulted in a flux of $(2.5 \pm 0.5) \gamma/(\text{cm}^2\text{s})$. A compatible result was found at Fermilab, with a measured flux of $(1.7 \pm 0.9) \gamma/(\text{cm}^2\text{s})$. Underground at INFN-LNGS, instead, the total flux was smaller, $(1.0 \pm 0.5) \gamma/(\text{cm}^2\text{s})$, due to the overall lower contamination of the surrounding rocks.

Fig.4.1b shows a typical energy spectrum, where the cut-off after the ^{208}Tl peak at 2.615 MeV is clearly visible. Notice that, in a qubit experiment, the chip is surrounded by the metal vessels of the dilution refrigerator and, occasionally, by an external mu-metal shield that protects it from external magnetic fields. These elements block a fraction of the incoming γ -rays, resulting in a reduction of the overall flux. The plot shows the energy spectrum in three different configurations: with the NaI spectrometer unprotected, with

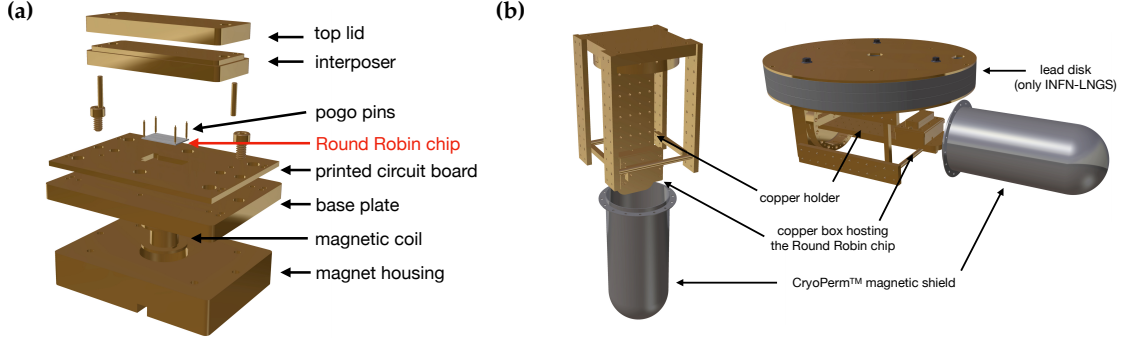


FIGURE 4.2: (a) Round Robin chip (gray) enclosed in its gold-plated copper box hosting the printed circuit board, and (b) its copper holder, with also the lead shield used at Gran Sasso Underground Laboratory for the mitigation of γ radioactivity.

the spectrometer surrounded by the metal vessels of the dilution refrigerator, and with it surrounded by the metal vessels of the dilution refrigerator plus the mu-metal shield.

4.2.2 Close sources

Close sources of radioactivity comprise every component that cannot be placed far away from the chip, like cables and electronic components. Qubit setups are not standardized, so different experiments will employ different components and in different number. Nonetheless, the vast majority of them are common to many experiments, so we can expect that different configurations would still result in similar contribution on qubit decays.

For this work, we considered the setup and the components used in the Round Robin experiment. The chip is a $325\ \mu\text{m}$ -thick, $11.9 \times 7\ \text{mm}^2$ high-resistivity silicon wafer hosting 16 transmon qubits (14 flux-tunable qubits and two fixed-frequency qubits) with the corresponding readout resonators. Since also interactions in the substrate are detrimental for qubit coherence, the entire chip, instead of only the qubits, was considered as a target in the simulation.

The chip is enclosed in a gold-plated copper box, above a printed circuit board (PCB) used for connecting it to the input and output RF cables (see Fig.4.2a). At the bottom of the box is hosted the magnetic coil used for tuning the transition frequencies of the qubits. The box is then placed on a copper holder, necessary both for hosting other electronic components and for granting thermal contact with the coldest state of the dilution refrigerator, and surrounded with a CryoPerm™ magnetic shield to reduce flux noise (see Fig.4.2b).

In Fig.4.2b it is possible to see also a 3 cm thick lead shield placed above the copper holder: this is part of the shielding against environmental γ -rays employed at the Gran

Sasso Underground Laboratory. This shielding has been taken into account also in the discussion of close sources because lead is contaminated with ^{210}Pb , a well-known γ emitter, with typical activity of about 100 Bq/kg.

Placed on the copper holder, or in close proximity with it, there are then all the other RF components necessary to send RF signals to the chip, so cables, attenuators, isolators, filters, and so on.

The full list of close sources considered in this work included:

- A* printed circuit board (PCB) – weight: ~ 7 grams. The PCBs used for the Round Robin experiment were produced by San Francisco Circuits. They consisted of a $1.575 \times 2.205 \text{ cm}^2$, 0.157 cm thick 3-layer FR408HR (a high performance laminate for multi-layer Printed Wiring Board applications, where maximum thermal performance is required). We also measured a new type of non-magnetic PCB, indicated with A^* in Table 4.1;
- B* gold-plated copper box (total mass: 0.4 kg) and a gold-plated copper holder (2.2 kg) used to thermally anchor the qubit to the mixing chamber plate as shown in Fig. 4.2;
- C* CryoPermTM magnetic shield, consisting in a hollow cylindrical shield with a diameter of 78.8 mm, 193.5 mm height and 1.0 mm thickness – weight: ~ 475 grams;
- D* Intelliconnect non-magnetic SMA adapters – weight: ~ 10 grams;
- E* Intelliconnect non-magnetic copper coaxial cables (diameter: 2.19 mm, length: 25 cm);
- F* Radiall SP6T cryogenic switch – weight: ~ 165 grams;
- G* 4–12 GHz Low Noise Factory circulator (mod. CIC4_12A) – weight: ~ 50 grams;
- H* 4–12 GHz Low Noise Factory dual-junction circulator (mod. CICIC4_12A) – weight: ~ 97 grams;
- I* 4–8 GHz Low Noise Factory triple-junction isolator (mod. ISISISC4_8A) – weight: ~ 145 grams;
- J* 10–30 dB XMA attenuators (mod. 2082-6418-10-CRYO, 2082-6418-20-CRYO, and 2082-6418-30-CRYO) – weight: ~ 5 grams each;
- K* K&L Low Pass Filters (mod. 3L250-12240/T20000 and 3L250-8160/T20000) – weight: ~ 15 grams each;
- L* COAX NbTi superconducting coaxial cables running from 4 K to 15 mK (mod SC-219/50-NbTi-NbTi).

We also considered other components placed at the 4 K stage of the dilution refrigerator, which are:

M 4–8 GHz Low Noise Factory cryogenic amplifier – weight: ~ 17 grams;

N COAX copper-beryllium coaxial cables running from 300 K to 4 K (mod SC-119/50-B-B).

For these sources, we expect their contributions on qubit decay rate to be smaller due to the higher distance from the chip.

Finally, we considered also two other materials widely used in qubit experiments, which are:

O Thermally conductive epoxy glue (StycastTM);

P DOW CORNING cryogenic grease.

Measurements of radioactive contamination were performed with High-Purity Germanium (HPGe) detectors, that allowed to reconstruct the energy spectrum of γ radiation from each source and, based on the heights of the peaks, the activities of the different contaminants.

Results are reported in Table 4.1. Of all the components, the PCB is by far the most contaminated one. This was expected, since the glass fiber used in their production is known for its high concentration of radioactive impurities [156–158]. We also considered a possible alternative type of a-magnetic PCB (A^*), consisting of copper (three layers, for a total thickness of $87\ \mu\text{m}$) interleaved by dielectric layers (hydrocarbon ceramic laminates) produced by ROGERS Corporation under the codes 4350B (two layers, for a total thickness of $420\ \mu\text{m}$) and 4450F (a single layer, $95\ \mu\text{m}$ thick). As shown in Table 4.1, these PCBs (A^*) were about a factor of 3 cleaner, from the radioactivity point of view, compared to those selected for the Round Robin chip.

On the contrary, the materials dominant in weight and closest to the chip - copper (B) and CryoPerm (C) - feature a good radiopurity level. Nevertheless, these components are the ones most subject to cosmogenic activation. Indeed, this analysis confirms the presence of cosmogenic radioisotopes - ^{54}Mn , ^{57}Co , ^{58}Co , and ^{60}Co - produced above-ground by cosmic ray interaction with the materials. The results are reported in Tab. 4.2.

Component	^{232}Th (mBq/kg)	^{238}U (mBq/kg)	^{235}U (mBq/kg)	^{40}K (mBq/kg)	^{137}Cs (mBq/kg)
<i>A</i>	(18000 $\pm$ 1000)	(11500 $\pm$ 400)	(710 $\pm$ 110)	(12000 $\pm$ 1000)	< 30
<i>A*</i>	(5410 $\pm$ 330)	(4200 $\pm$ 200)	(230 $\pm$ 50)	(4200 $\pm$ 500)	< 40
<i>B</i>	< 1.5	< 1.2	< 4	< 9	< 0.6
<i>C</i>	< 8.4	< 8.3	< 8.4	< 35	< 2.7
<i>D</i>	(46 $\pm$ 13)	(42 $\pm$ 10)	(70 $\pm$ 30)	(240 $\pm$ 90)	< 10
<i>E</i>	(54 $\pm$ 12)	(44 $\pm$ 11)	(34 $\pm$ 17)	(740 $\pm$ 130)	< 12
<i>F</i>	(1880 $\pm$ 100)	(1340 $\pm$ 60)	(130 $\pm$ 30)	(2200 $\pm$ 300)	< 11.2
<i>G</i>	< 310	< 330	< 410	< 2000	< 60
<i>H</i>	< 250	< 380	< 380	< 2600	< 60
<i>I</i>	< 190	< 240	< 220	< 2000	< 50
<i>J</i>	< 52	(200 $\pm$ 20)	< 47	< 140	< 13
<i>K</i>	(23 $\pm$ 4)	< 9.1	(60 $\pm$ 10)	< 100	< 1.9
<i>L</i>	< 750	< 1000	< 380	< 7000	< 230
<i>M</i>	< 890	< 1000	< 850	< 10000	< 210
<i>N</i>	(240 $\pm$ 40)	< 78	(350 $\pm$ 90)	< 500	< 20
<i>O</i>	(53 $\pm$ 4)	(9400 $\pm$ 900)	(350 $\pm$ 30)	(290 $\pm$ 40)	< 2.2
<i>P</i>	< 10	< 11	< 4.5	< 87	< 5

TABLE 4.1: Bulk contamination - values and 90% C.L. upper limits - of the materials located in the proximity of the qubit at the 15 mK stage (A - L) and 4K (M,N) stage. For the ^{232}Th decay chain, we quote the maximum activity between the ones measured for ^{228}Ra and ^{228}Th ; concerning ^{238}U and daughters, we quote the ^{226}Ra activity, being representative of the most worrisome part of the decay chain.

Component	^{60}Co (mBq/kg)	^{58}Co (mBq/kg)	^{57}Co (mBq/kg)	^{54}Mn (mBq/kg)
<i>B</i>	(0.6 ± 0.3)	(4 ± 1)	—	(2.4 ± 0.8)
<i>C</i>	< 3.7	(14 ± 5)	(20 ± 7)	—
<i>D</i>	(51 ± 8)	—	—	—
<i>K</i>	(5 ± 1)	—	—	—

TABLE 4.2: Activity concentration - values and 90% C.L. upper limits - for short- and long-lived radioisotopes produced by cosmogenic activation in the materials of the detector setup.

4.3 Estimating the impact

4.3.1 Geant4 simulation

When a particle crosses a material, its interactions with the surrounding atoms are governed by the cross sections of the physical processes that can occur. As a consequence, both its trajectory and energy exchanges are probabilistic. Two particles of the same species, entering the same material with identical energy and momentum, will end up following different paths due to the different outcomes of their interactions with the material atoms.

This makes Monte Carlo methods particularly useful when working with these processes. Random sampling is used to determine the outcome of a single interaction, and by repeatedly simulating the passage of particles through the volume of interest it is possible to reconstruct physical quantities like the spectrum of the deposited energy.

Geant4 simulations are based on this technique [146]. The trajectory of a particle is divided into small steps, and for each of them the interaction outcomes, such as energy and momentum exchange or the production of secondary particles, are determined through random sampling. This step-by-step process ends when the particle and its secondaries disappear due to an interaction (like decay or inelastic scattering), are slowed down to zero kinetic energy or reach the end of the simulation volume.

Fig.4.3 shows how the geometry of the qubit experiment was reconstructed in Geant4. Setup elements shown in Fig.4.2 are reconstructed from their CAD models, while for the dilution refrigerator a simplified geometry was used, including the plates corresponding to the different temperature stages and the cylindrical metal vessels. A cylindrical lead shield (visible in Fig. 4.4a) was also implemented in the simulation to estimate its impact on reducing the flux of incoming particles.

In our simulation, ionizing particles from far sources – γ -rays, muons and neutrons – are generated according to the momentum and energy distributions measured in the laboratory. The simulation allows us to estimate the fraction of these particles reaching

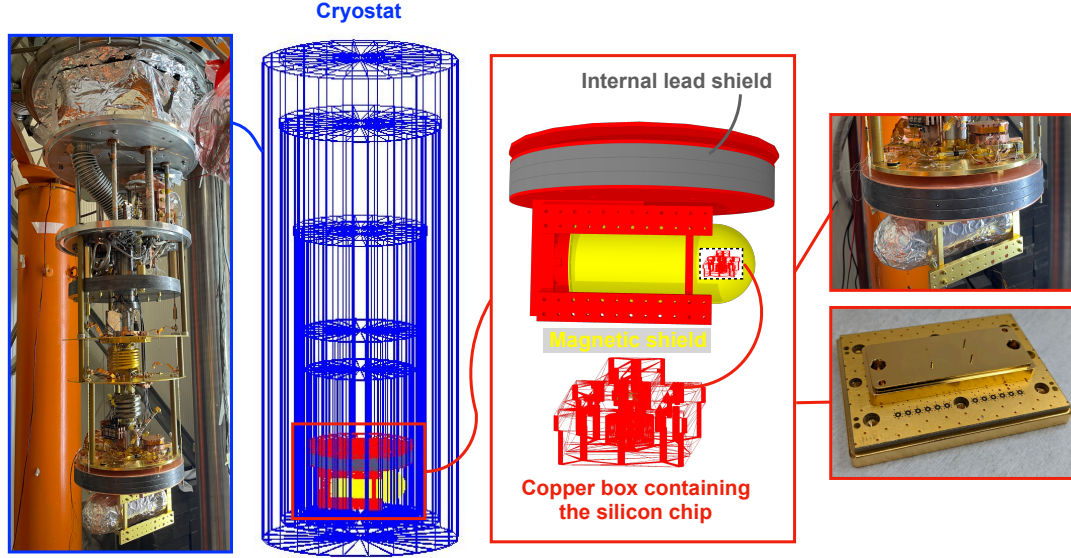


FIGURE 4.3: The experimental setup reconstructed in the Geant4 simulation.

the chip and depositing energy in it. From this fraction we can scale the measured flux and estimate the rate of impacts in the superconducting chip.

Neutrons and γ -rays were generated uniformly on the surface of a cylinder enclosing the cryostat, according to their measured energy spectra (see Fig.4.1) and with an isotropic momentum distribution. The greatest limitation in making this simulation is that only a very small fraction of the generated particles will actually reach the chip: for γ 's, we estimate that only one in 10^8 generated particles will actually produce an interaction in the chip. To reduce the statistical uncertainties in the rate of interactions and the spectrum of the deposited energy, then, it would be necessary to generate tens of billions of particles, which would require excessively long computational times.

To overcome this limitation we divided our simulation into two steps (see Fig.4.4). First, we generated particles uniformly on a cylinder S1 (Fig.4.4a) surrounding the entire cryostat with its shielding. Instead of directly measuring the fraction of them reaching the chip, though, we measured the fraction of them reaching a smaller cylinder S2 surrounding the setup, along with their energy spectrum and momentum distribution (Fig.4.4b and Fig.4.4c). We then generated γ 's on the S2 surface according to the energy and momentum distribution previously obtained, and measured the fraction of them depositing energy in the chip. Finally, to obtain the rate of impacts in the chip, we divided the number of events depositing energy in the chip by the equivalent time of the simulation, which is given by $t_{eq} = N_{gen}/(A_2 \cdot \text{flux}_2)$, where N_{gen} is the number of generated events, A_2 is the area of the cylindrical surface used to generate gammas in the second step, and flux_2 is the gamma-ray flux calculated in S2 after the first step.

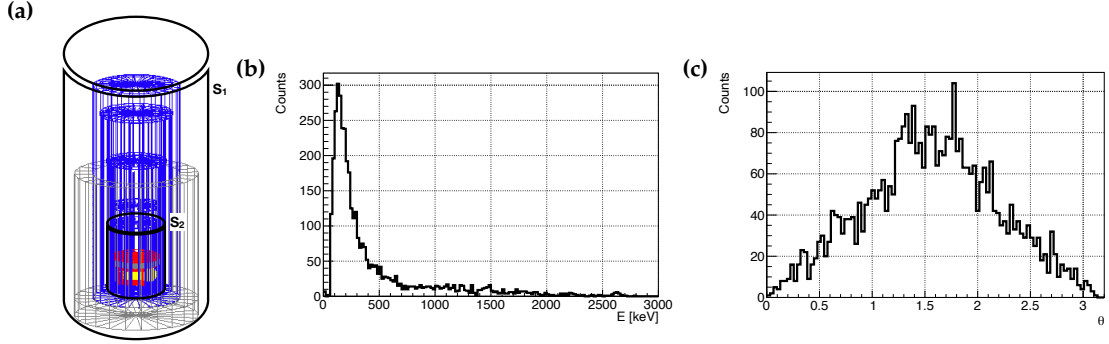


FIGURE 4.4: (a) Experimental setup with both the internal and the external lead shield, along with the surfaces S1 and S2 used for the double-step simulation of the interactions of γ -rays. (b) Energy distribution of γ 's entering S2 and (c) their polar angle distribution.

Cosmic muons, instead, do not come from every direction, since their zenith angle follows a $\cos^2$ distribution¹. Because of that, their generation follows a different approach. First, positions on a hemisphere surrounding the cryostat are generated according to the $\cos^2$ distribution. Then, for every sampled position, muons are generated perpendicular to a $(120 \times 120) \text{ cm}^2$ plane tangent to the hemisphere at the position previously generated.

Finally, close sources were simulated by generating radioactive decays of relevant isotopes uniformly distributed within the volume of each setup component (components A to N of Table 4.1). The location of every component in the simulated geometry was assumed to be the typical location where the given item is usually mounted in the cryostat. To estimate the rate of impacts in the Round Robin chip, for each setup component we scaled the measured activity by the fraction of decays producing interactions in the chip, obtained from the simulation.

4.3.2 Simulation results

For each far radiation source, the corresponding rate of particle impacts produced in the chip was estimated both for a “standard” above-ground laboratory, using average values for the fluxes of far sources, and deep-underground at the Gran Sasso Underground Laboratory (INFN-LNGS). In the case of LNGS, three different levels of shielding were considered: no shield at all, a 10-cm-thick lead shield surrounding the dilution refrigerator, and the so-called “full” shield configuration, with both the external lead shield and the 3-cm-thick inner lead shield placed right above the copper holder of the setup (see Fig.4.2b).

¹In deep-underground laboratories the rock overburden affects the angular distribution, but this does not impact significantly the simulation result.

Location	Rate (mHz)		
	Lab. γ -rays	Muons	Neutrons
“Standard”	(18 ± 4)	(10.0 ± 0.6)	(0.15 ± 0.05)
LNGS	(7.0 ± 3.5)	$< 10^{-5}$	$< 10^{-4}$
LNGS Ext. Shield	(0.10 ± 0.05)	$< 10^{-5}$	$< 10^{-4}$
LNGS Full Shield	(0.07 ± 0.04)	$< 10^{-5}$	$< 10^{-4}$

TABLE 4.3: Interaction rate in the substrate of the Round Robin chip. For the above-ground facilities we assumed a “standard” γ -ray flux of $(2.5 \pm 0.5) \gamma/\text{cm}^2/\text{sec}$ and a muon flux of $1 \mu/\text{cm}^2/\text{min}$. For the LNGS facility we used the measured flux of $(1.0 \pm 0.5) \gamma/\text{cm}^2/\text{sec}$ and we evaluated the suppression factor due to the presence of lead shields.

Component	Description	Rate (mHz)
<i>A</i>	PCB	4.52 ± 0.04
<i>B</i>	Box	$[1 - 6] \times 10^{-3}$
<i>B</i> *	Holder	$[2 - 4] \times 10^{-4}$
<i>C</i>	Magnetic Shield	$[2 - 9] \times 10^{-4}$
<i>D</i>	SMA	$(2 \pm 0.4) \times 10^{-5}$
<i>E</i>	Cu coax cables	$(3 \pm 0.6) \times 10^{-5}$
<i>F</i>	Cryogenic switch	$(1.0 \pm 0.2) \times 10^{-2}$
<i>G</i>	Circulator	$< 8 \times 10^{-4}$
<i>H</i>	Dual-junct. circulator	$< 2 \times 10^{-3}$
<i>I</i>	Triple-junct. isolator	$< 2 \times 10^{-3}$
<i>J</i>	Attenuators	$[0.5 - 1] \times 10^{-5}$
<i>K</i>	Low Pass Filters	$(1 \pm 0.2) \times 10^{-5}$
<i>L</i>	NbTi cables	$< 4 \times 10^{-4}$
<i>M</i>	Cryogenic amplifier	$< 2 \times 10^{-5}$
<i>N</i>	Cu-Be cables	1×10^{-6}

TABLE 4.4: Interaction rate in the substrate of the Round Robin chip. The copper box and holder were split in this Table to highlight that, due to its proximity to the qubit, the contribution of the box is more important than the one of the holder, despite the same level of radiopurity. Some components are highlighted in bold to remind that usually these components are placed in large quantity in the dilution refrigerator (and the rate should be scaled accordingly).

Simulation results for far sources are summarized in Table 4.3. The greatest contributions come from environmental γ -rays, for which we estimated a rate of impacts in the chip of (18 ± 4) mHz. At INFN-LNGS, instead, the concentration of radioisotopes in the Gran Sasso rock is lower, resulting in a rate, without any type of shielding, of only (7.0 ± 3.5) mHz. This can be further reduced by implementing the lead shielding: an ideal cylinder, 8 cm tall and 10 cm thick, can suppress this rate by a factor 70, down to (0.10 ± 0.05) mHz; the 3-cm-thick lead shield placed right above the copper holder is capable of further reducing the rate of impacts down to (0.07 ± 0.04) mHz. Notice, however, that lead shields are rarely ideal cylinders: more commonly, they are made out of lead bricks, resulting in a less tight geometry. This estimate is therefore meant to provide only an order-of-magnitude indication of the achievable suppression. In Chapter 6, we will present measurements performed on a transmon device operated with a realistic external lead shield; in that case, the corresponding reduction factor was estimated to be around 35.

As mentioned in section 4.2.2, shielding the experiment with lead could in principle also be problematic, due to the presence of ^{210}Pb among its isotopes. This was concerning especially for the internal lead shield, which is the closest one to the chip. However, when estimating its contribution we obtained a rate of impacts in the Round Robin chip of only 0.01 mHz, an increase that is negligible compared to the amount of radiation that it is blocking.

Table 4.3 also reports the contributions from cosmic muons and neutrons. Muons account for the second most significant contribution in above-ground laboratories, with a rate of impacts in the chip of (10 ± 0.6) mHz. Despite their rate being lower than the one of environmental γ -rays, their presence can still be worrisome for quantum computers: as already observed by Wilen et al. [131], muons, when crossing the substrate, leave behind a long track of electrons and holes, longer than the one of an electron being kicked by the Compton scattering with a γ -ray. As a consequence, muon interactions can affect more qubits on the same chip, complicating further the implementation of quantum error correction.

In the case of muons, shielding them is also more complicated compared to environmental γ -rays. For them, the only possible shielding consists in moving the experiment to an underground site: at INFN-LNGS, their impact rate is reduced by six orders of magnitude, down to a negligible value ($< 10^{-5}$ mHz), but also facilities not as deep are suitable. As an example, the NEXUS (Northwestern Experimental Underground Site) underground laboratory at Fermilab, Illinois, US, is shielded by only 100 m of soil/gravel, and yet offers a suppression by three orders of magnitude of atmospheric muons [159, 160].

An alternative to moving to an underground facility could be tagging incoming muons with an external detector in order to identify and reject operations done while particles are traversing the chip, a proposal that was shortly discussed in Section 3.3.3. Placing a big muon detector out of the cryostat, though, would result in a rate of triggers as high as hundreds of Hz [161]. A more efficient veto system could be made out of cryogenic detectors in the proximity of the chip, that would reduce the trigger rate down to the only ones actually impacting qubits [138, 162], like the one tested in our group and described in Appendix A.

There are also proposals [163] to minimize the impact of muon interactions by re-designing quantum error correction algorithms: the idea is to distribute the qubits employed across several chips, so when one of the chip gets hit by a muon, this does not necessarily translate into a general failure of the algorithm.

Concerning neutrons, which are the main concern for classical computers [164], we found that their contribution on qubit decays is actually much lower than the one of other far sources: their impact rate in an above-ground laboratory is just (0.15 ± 0.05) mHz, reducing to negligible values ($< 10^{-4}$ mHz) when moving the experiment deep underground.

Table 4.4 summarizes, instead, the simulation results for close sources. For all of them, the vast majority of the interactions are coming from γ -rays emitted during decays, with α and β particles playing a role only when emitted by the PCB or the copper box, which are the setup elements closest to the chip.

Of all contributions, the most worrisome is by far that from the PCB, which was estimated to be (4.52 ± 0.04) mHz. Other sources produced impact rates negligible compared to that one: even the contribution from the cryogenic switch, which is the second highest one in Table 4.4, is two orders of magnitude lower than the contribution from the PCB.

There are two reasons why the PCB is so problematic: it is, along with the copper box, the closest element to the chip; its level of radiopurity, also, is poor compared to other materials. This problem is well known to particle physicists developing detectors for rare event searches, as both fiberglass and ceramic (typical materials used for the PCB multilayers) usually contain levels of radioactivity comparable to those in rocks and soil. When comparing its contribution with the ones from far sources, we can see that deep underground it becomes the main source of interactions if the setup is shielded against external γ radiation. In above-ground laboratories, instead, its contribution is lower than the ones from environmental γ -rays and cosmic muons, but still accounts for more than 10% of the total rate of impacts. This poses a limit on how much the radioactivity level can be reduced by employing only shielding techniques.

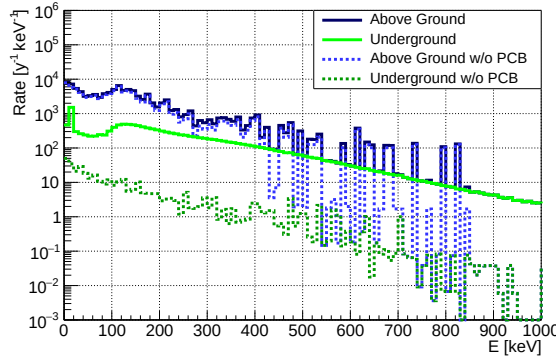


FIGURE 4.5: Rate of interactions in the Round Robin chip as a function of the deposited energy. In blue is the result for an above-ground experiment, while in green is the same result for an experiment at INFN-LNGS. Dashed lines are the results without the contribution from the PCB.

This limit can be improved by employing new materials for the fabrication of the PCB. In our research we investigated also the contribution coming from a novel non-magnetic PCB with better levels of radiopurity (reported as A^* in Table 4.1). The simulation showed that the rate of impacts from this new PCB is approximately a factor of 3 lower than the one coming from the PCB designed for the Round Robin experiment, confirming that progresses on this aspect are possible.

Summing close and far sources, we found an overall rate of impacts in the Round Robin chip of approximately 33 mHz, around one event every 30 seconds. It's the same order of magnitude found by Wilen et al. [131] in their research. Figure 4.5 shows the rate as a function of the energy deposited in the chip. The majority of interactions deposit tens/hundreds of keV, with highest deposits coming from decays in the PCB.

In conclusion, we proved that far sources, especially environmental γ -rays and cosmic muons, are responsible for the majority of the particle impacts on qubit chips. Lead shielding can help abating the contribution from γ particles, but to mitigate also the rate of impact from muons it is necessary to move the experiments to underground facilities. Having mitigated the contributions from far sources, the contribution from the PCB becomes the most significant one, accounting for one impact every few minutes. Abatement of this contribution, though, is possible only through dedicated R&D activities aiming at the employment of different materials for the PCB fabrication.

Chapter 5

Impact of radioactivity on the decay time of a fluxonium qubit

The study described in the previous chapter showed that far sources, in particular environmental γ -rays and cosmic muons, are the main contributors to particle interactions in quantum devices. Our next step was to estimate the extent to which these sources affect the qubit performance.

Estimating their impact, though, is not trivial: as described in Section 2.3, qubit decoherence is the consequence of many different phenomena, and disentangling radioactivity from the others can be difficult. For this reason, we adopted a comparative approach: our chips are first characterized in a shielded configuration, deep underground to mitigate interactions from cosmic muons and surrounded by a lead shield to stop environmental γ -rays; subsequently, the measurement is repeated while exposing the chip to ionizing radiation. The specific contribution from radioactivity is then inferred by comparing the results of the two measurements.

The first device we employed for our experiment was a fluxonium qubit designed and produced at the Karlsruhe Institute of Technology (KIT), in Germany. Measurements were performed in the Ieti Underground Facility at INFN-LNGS, where the flux of cosmic muons is attenuated by a factor 10^6 . The qubit was characterized in three different configurations: with an external lead shield surrounding the dilution refrigerator, in order to block environmental γ -rays; without the external lead shield; without the external lead shield but with the addition of a γ radiation source placed next to the cryostat.

In the experiment we focused on the decay time of the qubit, looking for fluctuations correlated with the exposure to ionizing radiation. We then analyzed the time distribution of the spontaneous jumps of the qubit from $|0\rangle$ to $|1\rangle$ and vice versa, looking for deviations from its typical behavior caused by radiation-induced quasiparticle poisoning.

5.1 Experimental setup

A scheme of the experimental setup is depicted in Fig. 5.1. The qubit is placed at the coldest stage of the dilution refrigerator, where the base temperature is approximately 10 mK. The device employed works in reflection, so with a single RF connector used both for input and output. A second connector was used to send drive pulses from a different line. In order to prevent the reflected signal from interfering with the next incoming signals, a circulator was added to the line, sending the incoming readout signal to the chip and the reflected signal to the amplifiers.

The RF lines for readout and drive are designed following the same idea described in section 2.2. Readout and drive signals are generated at room temperature with much higher amplitude than the one required by the qubit, in order to make the contribution from thermal noise negligible. They are then attenuated inside the dilution refrigerator, where the thermal noise is lower, before reaching the qubit. Following the same principle, outgoing signals are amplified in multiple stages already inside the dilution refrigerator, before being demodulated at room temperature.

Pulses are generated by an FPGA-based board and up-converted to the frequency of the readout resonator by mixing with the signal coming from a local oscillator (LO). They are then attenuated by 60 dB before reaching the device, with a 30 dB attenuator at the 4 K stage and two other attenuators, of 20 dB and 10 dB respectively, at the 10 mK stage. Note that the RF line itself provides an attenuation of approximately 7 dB at the frequency of interest, so a signal entering the cryostat, reaching the 10 mK stage and then going out again is attenuated overall by around 15 dB even if there are no other elements on the lines.

Thermal noise, however, is not the only noise source to be taken into account when preparing a qubit experiment. As discussed in section 2.3, infrared photons produced by blackbody radiation at higher temperature stages can also be detrimental for qubit coherence, being a well-known quasiparticle source. To prevent it from reaching the qubit, the chip is shielded inside a copper box (see Fig. 5.2b), while IR photons traveling through the RF lines are blocked by low-pass filters. In particular, two IR filters are

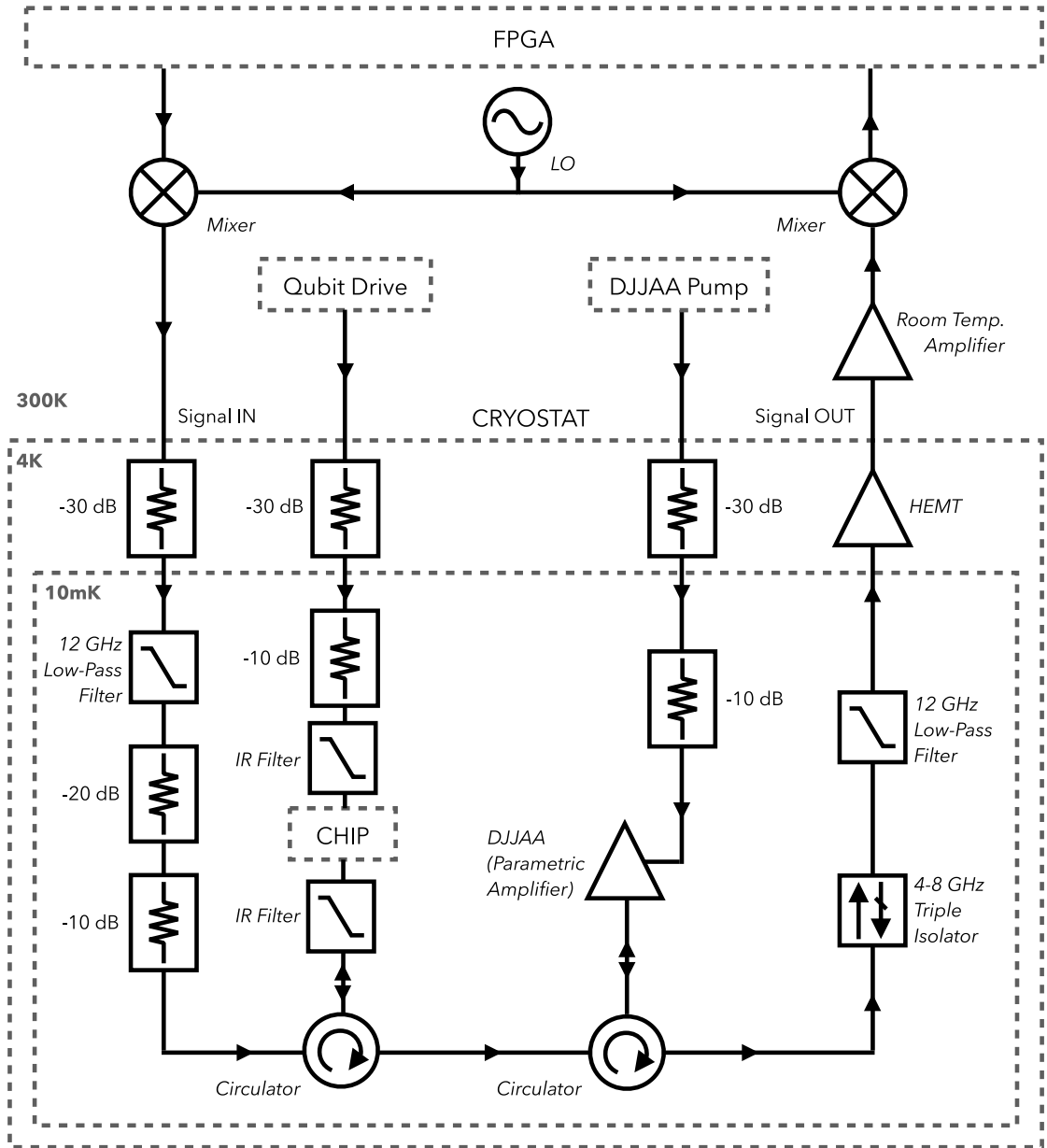


FIGURE 5.1: Scheme of the setup used for the experiment.

placed next to the chip, at the ends of both its ports, and two additional 12 GHz low-pass filters were added on the readout line, one on the input, right after reaching the 10 mK stage, and the other at the output, before leaving the 10 mK stage.

Amplification of the output signal was achieved in three stages. At 10 mK there was a parametric amplifier, the Dimer Josephson-Junction-Array Amplifier (DJJAA) [165], that will be described in section 5.1.2. As for the chip, also the DJJAA works in reflection, so it was connected to the readout line through a circulator to prevent amplified signals from going back to the qubit and affect the system.

The second amplifier was a High Electron Mobility Transistor (HEMT) Amplifier placed

at the 4 K stage. To prevent noise from the HEMT from reaching the qubit, a triple-junction isolator was also added on the output line, at the 10 mK stage. Following the HEMT, a third amplifier was placed at room temperature, for a total amplification of approximately 90 dB.

Finally, the amplified signal was down-converted to its original frequency by mixing with the same LO signal used in the up-conversion, and sent to the FPGA for the readout.

Since both the qubit and the parametric amplifier can be tuned by means of an external magnetic flux, two DC lines were also implemented, connecting two DC signal generators at room temperature to the coils placed next to the devices.

5.1.1 The qubit

The qubit employed is a fluxonium qubit produced at the Karlsruhe Institute of Technology (KIT). It is fabricated on a $15 \times 18 \times 0.33 \text{ mm}^3$ c-plane sapphire substrate, and its peculiarity is that the shunting inductance, instead of being obtained with a series of tens of Josephson junctions, is made of a granular aluminum (grAl) superconducting strip [63].

grAl films are obtained when depositing aluminum on a substrate in a controlled oxygen atmosphere: the result is a mixture of pure aluminum grains of nanometric dimension and aluminum oxide. This is conceptually similar to what is done in a standard fluxonium qubit, where a superconducting strip is interrupted by tens of Josephson junctions in order to reach high inductance values. The coupling between the grAl grains determines the normal-state resistivity of the material and its critical temperature, and can be varied during fabrication by properly tuning the oxygen pressure, the deposition rate and the substrate temperature [166]. The kinetic inductance of the superconductor is proportional to the normal-state resistivity: the result is that, when the grAl strip is produced with the right coupling between grains, the mixture of aluminum oxide and pure aluminum grains features an inductance comparable with that achieved for the shunting inductance of other fluxonium qubits.

A scheme of the qubit with its readout resonator is shown in Fig. 5.2a: the qubit consists of a SQUID shunted by an inductance L and coupled to the readout resonator through a shared inductance L_s . As already discussed in the last section, two RF lines are used for the readout and for the drive pulse, respectively.

Fig. 5.2b shows a photo of the qubit inside the copper box used to shield it against infrared radiation. The box is attached to a copper rod used for thermalization. In the

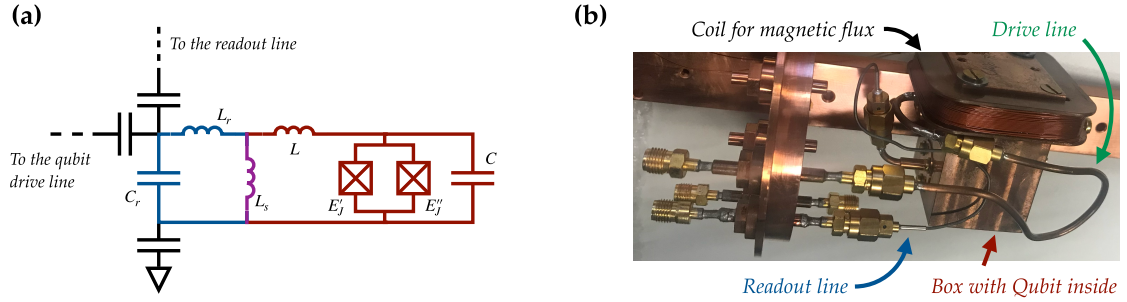


FIGURE 5.2: (a) Scheme of the fluxonium qubit employed in the experiment. C_r and L_r are, respectively, the inductor and the capacitor of the readout resonator, inductively coupled to the qubit by L_s . L is the shunting inductance of the fluxonium, which also features two Josephson junctions. The capacitor C takes into account capacitive effects of the junctions. (b) Photo of the qubit inside its copper box, along with the coil used for magnetic flux biasing and the connectors for the readout and drive lines.

photo it is also possible to see the coil used to generate the magnetic flux necessary for biasing and the readout and drive lines exiting the copper box.

5.1.2 The DJJAA

The Dimer Josephson-Junction-Array Amplifier (DJJAA) employed in our experiment is shown in Fig. 5.3a. Its circuit consists of an array of thousands of identical SQUIDS in series, interrupted in the middle by a large capacitance C_c (see Fig. 5.3b). Each SQUID is connected to ground by means of a capacitance C_0 . The array is connected, on one end, to a transmission line used both for input and output, while the other end is terminated to ground.

The working principle is similar to the one described in section 2.2 for Josephson Parametric Amplifiers, with the readout signal being amplified by draining energy from a *pump* signal (in Fig. 5.3a it is possible to see the RF connectors for the signal and pump). There are some differences, though. JPAs, in fact, amplify signals only at one resonant frequency, or *mode*; by replacing the single Josephson junction with an array of hundreds, or thousands, of junctions, the number of frequency modes in the range of interest, for which the signal is amplified, increases. If, instead of single junctions, the array is made of SQUIDS, like in the DJJAA, the system also becomes flux-tunable (the coil can also be seen in Fig. 5.3a). The result is an amplifier with a range spanning several GHz, since different frequencies can be reached by properly choosing the resonant mode and the magnetic flux biasing.

In order to understand the reason why the array of SQUIDS is interrupted by the central capacitance, it is necessary to consider another limitation of the toy model of parametric amplifier described so far: the amplification is dependent on the phase of the pump

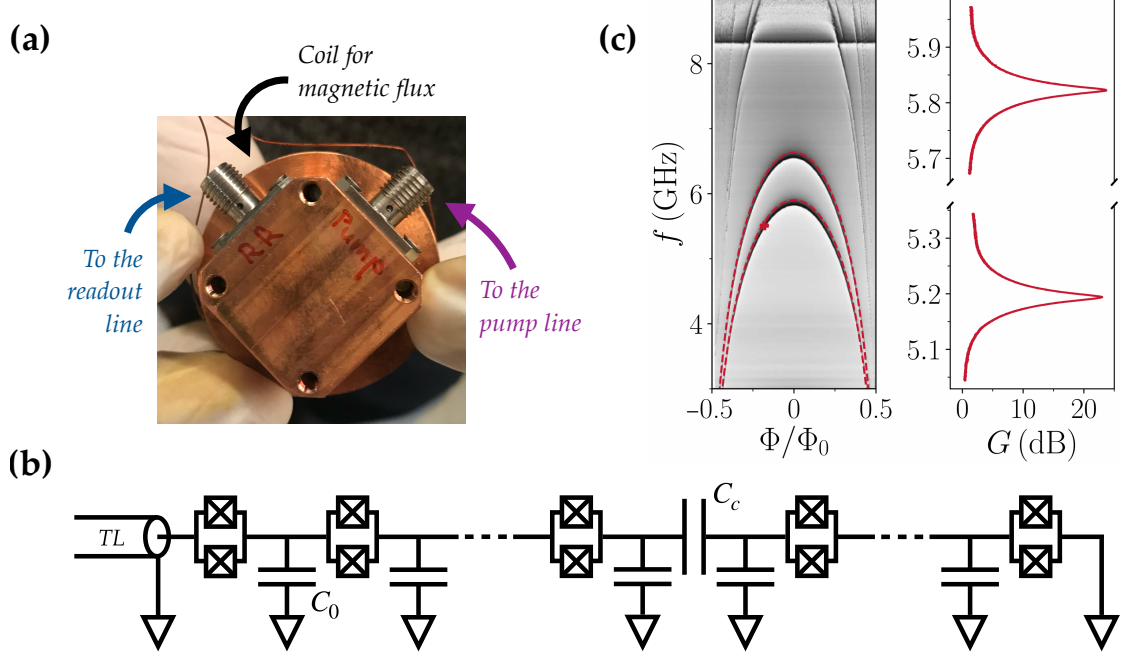


FIGURE 5.3: (a) Photo of the DJJAA employed for the experiment. (b) Scheme of the DJJAA. (c) On the left, S_{21} signal from the DJJAA as a function of the magnetic flux bias; on the right, its gain curve; both plots are taken from [156]

signal. If the phase difference between the pump and the signal is not the correct one, the signal could even be de-amplified. This type of amplifier is also known as *degenerate parametric amplifier* [75], because of the relation between the signal mode and an *idler* mode. Due to energy conservation the energy drained from the pump cannot go entirely to the signal, but has to be redistributed according to one of the two relations:

$$\omega_p = \omega_s + \omega_i \quad (5.1a)$$

$$2\omega_p = \omega_s + \omega_i \quad (5.1b)$$

Where ω_p is the pump frequency, ω_s is the signal frequency and ω_i is known as the idler frequency. The mixing process, i.e. whether the idler and the signal photons mix with one or two pump photons, depends on the nature of the non-linear coupling between the fields. In phase-sensitive parametric amplifiers, the idler mode and the signal mode correspond to the same frequency mode, and that is the reason why this type of amplification has also been named *degenerate*.

To avoid phase-sensitive amplification it is necessary to employ *non-degenerate parametric amplifiers*, where the idler mode and the signal mode have different frequencies.

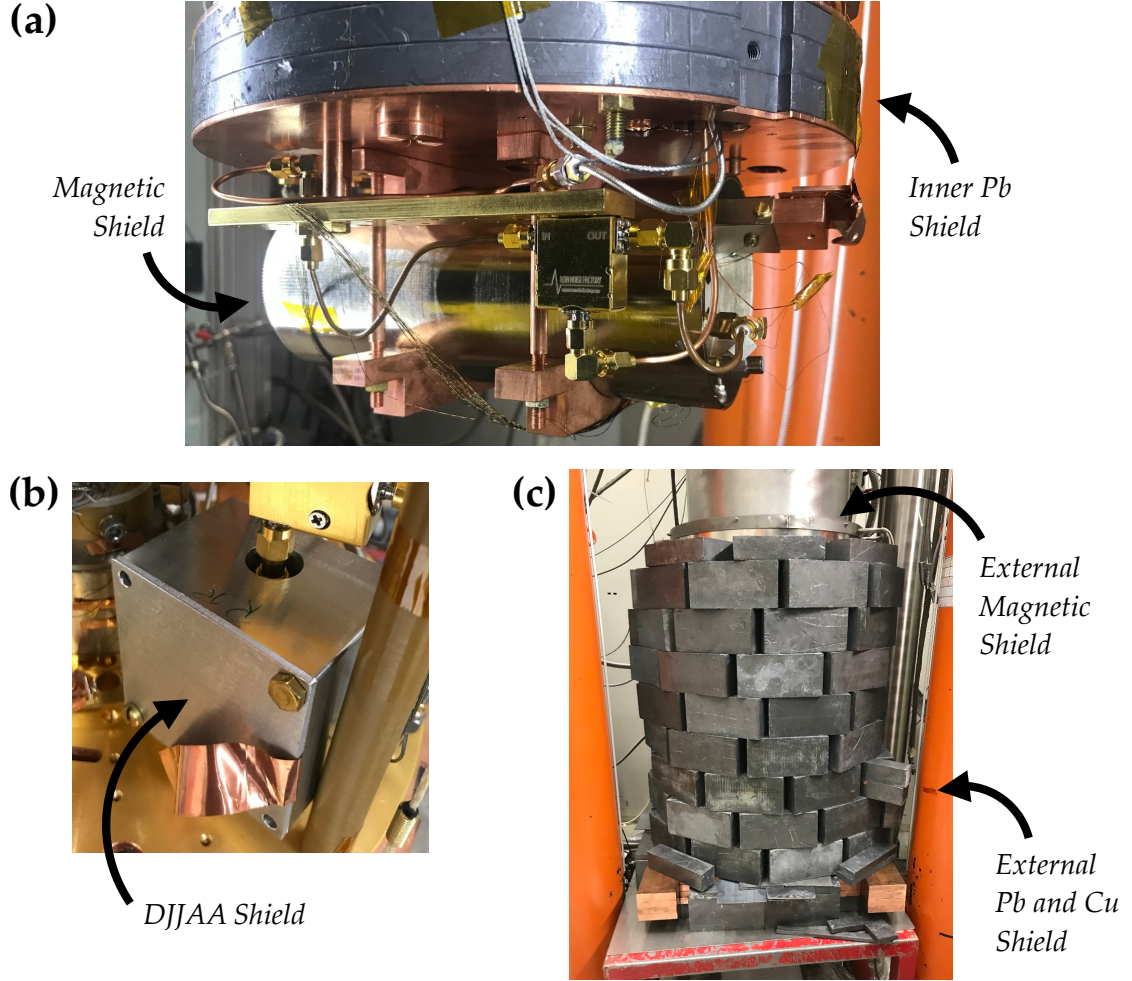


FIGURE 5.4: Photos of the shielding; (a) the magnetic shield surrounding the qubit, with the inner lead shield above it; (b) the shield for the DJJAA; (c) the external magnetic shield, out of the cryostat, surrounded with the lead and copper shield.

In the DJJAA, non-degenerate parametric amplification is achieved by adding a capacitance C_c in the middle of the SQUID array (see Fig. 5.3b). In this way, each mode hybridizes into a pair of coupled modes, i.e. a *dimer*, corresponding to symmetric and antisymmetric combinations of the two array sections, with the pump frequency in the middle, at the same distance from both modes.

The response of this type of amplifier is depicted in Fig. 5.3c. The left panel shows the first dimer, with the frequencies of idler and signal varying with the external magnetic flux Φ (there expressed in terms of the flux quantum Φ_0). At higher frequencies, other modes are also observable. The right panel shows the gains corresponding to the two modes when the pump is properly tuned.

Lead On [events/sec]	Lead Off [events/sec]	Source On [events/sec]
$(0.40 \pm 0.09) \times 10^{-3}$	$(19 \pm 1) \times 10^{-3}$	1.3 ± 0.2

TABLE 5.1: Expected rates of interactions in the chip in the three different configurations, estimated with simulations.

5.1.3 The shielding

The shielding employed in the experiment can be seen in Fig. 5.4. The qubit is boxed into a mu-metal shield to protect it against external magnetic fields, with a 3-cm-thick lead shield right above it to block environmental γ -rays coming from the top (Fig. 5.4a). A second 3-cm-thick lead shield, not visible in Fig. 5.4, was placed at the 4K stage. Since also the DJJAA is sensitive to external magnetic fields, it also has a small shield surrounding it, as can be seen in Fig. 5.4b.

The final layer of shielding was instead placed out of the cryostat, and can be seen in Fig. 5.4c. First, there was a two-layer Permimphy shield, to grant better protection against external magnetic fields. Then, there was a 5-cm-thick lead shield, made of $5 \times 10 \times 20$ cm bricks, to stop environmental γ -rays coming from the sides and from the bottom of the cryostat. In between the lead and the Permimphy, 1-cm-thick copper rods were also added to stop bremsstrahlung photons produced by interactions inside lead. The whole facility was set in the Hall C of INFN-LNGS, providing a shield also against cosmic muons.

The experiment was conducted with three different radiation exposures: with the full shielding against environmental radioactivity; without the external lead and copper shield (so only with the lead shields placed inside the cryostat); and with a Thorium source, placed next to the dilution refrigerator in order to increase the rate of radiation impacts in the chip. The source did not need to be placed inside the cryostat because most isotopes in the Thorium decay chain are emitting γ radiation, which can easily penetrate the metal vessels in the absence of heavy shielding. The total activity of the source was 161 kBq.

To estimate radiation contribution in these three configurations, we adapted the simulation framework previously developed for the Round Robin experiment and described in Chapter 4. Results of the simulations are reported in Table 5.1. The external shielding alone mitigates the total rate of interactions in the chip by more than one order of magnitude. Since this chip does not employ a PCB, which was the major setup contribution for the Round Robin experiment, interactions underground are due almost entirely to environmental γ -rays, with negligible contributions from the other sources.

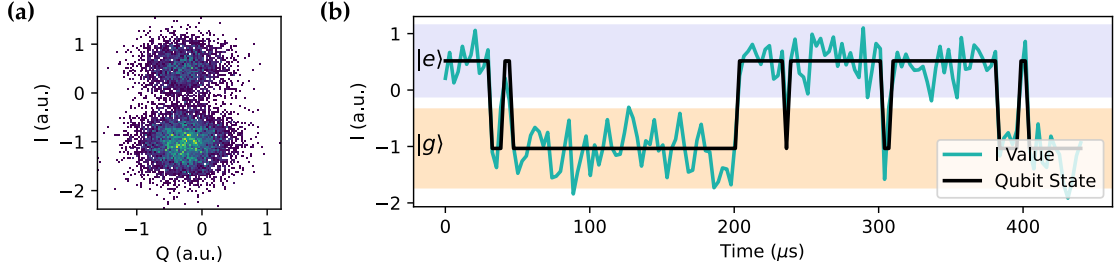


FIGURE 5.5: (a) Histogram of the IQ values of the demodulated readout signals: it is possible to identify two 2D Gaussian distributions, the upper one corresponding to the qubit being in $|e\rangle$ and the lower one to the qubit being in $|g\rangle$; the two distributions are aligned along the y axis, so the value of I alone is enough for state discrimination. (b) I measurements over time, with the corresponding qubit state; the two colored areas correspond to 2σ ranges from the central value.

A great contribution comes, instead, from the Thorium source, that increases the rate of interactions in the chip up to (1.3 ± 0.2) events/sec.

5.2 Measuring the energy-relaxation time

5.2.1 Measurement strategy

To investigate the effect of radioactivity on the fluxonium qubit, we chose the decay time T_1 as our estimator, following the approach of Vepsäläinen et al. [118] (see Section 3.1). Its measurement, though, followed a different method than the one described in section 2.3.3. For this experiment, we decided to measure the rate of spontaneous $|0\rangle \rightarrow |1\rangle$ and $|1\rangle \rightarrow |0\rangle$ quantum jumps of the qubit by repeatedly sending readout pulses, acquiring a state measurement approximately every $3 \mu\text{s}$.

State measurement is achieved through dispersive shift, as described in section 2.2: I and Q values of the demodulated signals form a double-Gaussian distribution on the $I - Q$ plane, with each Gaussian corresponding to a different qubit state. To simplify the analysis, I and Q are rotated in such a way that the separation between the two clouds is maximum along the I axis, making the information about Q not necessary (see Fig. 5.5a).

To determine whether the qubit spontaneously jumped between the two states, it was chosen to employ a *latching filter*, already employed in the analysis of quantum jump measurements in fluxonium qubits [167, 168]. The idea is that, to declare a jump, the I value of a qubit measurement has to lie within 2σ from the average I value of the corresponding final state (see Fig. 5.5b).

Once the $|0\rangle \rightarrow |1\rangle$ and $|1\rangle \rightarrow |0\rangle$ jumps were identified, it was possible to determine the transition rates $\Gamma_{\uparrow}$ and $\Gamma_{\downarrow}$. This was done by considering a stream of 10^6 measurements, so approximately 3 s, and by determining, in that interval, the times $T_{|0\rangle}$ and $T_{|1\rangle}$ the qubit spent in $|0\rangle$ and $|1\rangle$, respectively. Given $n_{\uparrow}$ and $n_{\downarrow}$ the number of $|0\rangle \rightarrow |1\rangle$ and $|1\rangle \rightarrow |0\rangle$ jumps, respectively, the two transition rates are computed as:

$$\Gamma_{\downarrow} = \frac{n_{\downarrow}}{T_{|1\rangle}} \quad \Gamma_{\uparrow} = \frac{n_{\uparrow}}{T_{|0\rangle}} \quad (5.2)$$

So the number of $|1\rangle \rightarrow |0\rangle$ jumps divided by the time the qubit spent in $|1\rangle$ and the number of $|0\rangle \rightarrow |1\rangle$ jumps divided by the time the qubit spent in $|0\rangle$. The energy-relaxation time is then determined by:

$$T_1 = \frac{1}{\Gamma_{\uparrow} + \Gamma_{\downarrow}} \quad (5.3)$$

Determining both $\Gamma_{\uparrow}$ and $\Gamma_{\downarrow}$ is particularly useful when looking for radiation-induced quasiparticle poisoning, since previous experiments have shown that, in the presence of quasiparticles, the distribution of time separation between consecutive jumps deviates from the exponential distribution expected by a random, Poissonian process [169].

In our experiment, we then looked for two different signatures of ionizing radiation: a statistically significant variation of T_1 when comparing the shielded and the unshielded configuration, that will be the topic of sections 5.2.2 and 5.2.3, and a variation in the jump statistics that can be ascribed to quasiparticle poisoning, that will be instead the topic of section 5.3.

5.2.2 “Long” measurements

In order to better estimate both the average T_1 value and its standard deviation, we decided that for each configuration we would have measured for approximately 7 hours. Given that each measurement cycle lasts a few seconds, this corresponds to $O(10^3)$ values of T_1 for each configuration.

The result of our measurements is plotted in Fig. 5.6. The left panel shows a plot with the superimposed histograms of the T_1 values in the three different configurations. On the right, instead, there are three plots, one per configuration, depicting the values of T_1 over time. In the T_1 vs time plot for the fully shielded configuration (red line) there are a few hours with no data: this was due to an error in the software for the control of the

	Lead On	Lead Off	Lead Off, Source On
Average T_1 [μs]	18.8 ± 0.1	19.8 ± 0.1	17.7 ± 0.1

TABLE 5.2: Average T_1 values, in μs , for the three configurations in the long measurements.

readout electronics, that stopped the acquisition. The total amount of measurements for that dataset, however, is still comparable with the other two.

The average T_1 values for the three configurations are reported in Table 5.2. They range from $(17.7 \pm 0.1) \mu\text{s}$, when the qubit is exposed to the Thorium source, to $(19.8 \pm 0.1) \mu\text{s}$ when the external shielding has been removed but the Thorium source is not present. These values do not correlate well with the radiation exposure: the average T_1 is lower when the qubit is fully shielded than when the external shield has been removed, meaning that this variation is due to different decoherence sources. However, when the qubit is exposed to the Thorium source there is a “tail” of low T_1 values. This could point to a strong effect of the source on qubit decoherence, but there are a couple of considerations actually disproving this conclusion.

First of all, a similar tail is not observed in the configuration without the external lead shield (the actual amount of low- T_1 measurements is slightly lower in this configuration than in the fully shielded one). It is true that the Thorium source produces a rate of particle interactions in the qubit chip much higher than environmental radioactivity, so we expect its impact on qubit coherence to be stronger. However, as seen in Section 5.1.3, the external lead shield can severely abate the contribution from environmental γ -rays: not observing a difference between the shielded and the unshielded configuration, but only when the Thorium source is present, weakens the hypothesis that this difference is radiation-induced.

Second, by looking at the time evolution of T_1 when the qubit was exposed to the source (bottom right plot, green line), it is possible to see that the majority of the low- T_1 measurements are concentrated in the first two hours of acquisition. If these values were radiation-induced, we would have expected to observe them in the entire dataset, since the activity of the source, and hence the rate of impacts in the chip, can be considered to be constant over time. We can also see that a similar “noisy” regime of the qubit was observed, even if not as strong, also in the fully shielded configuration: this happened in the first two hours and then in the last one. This further confirms that different decoherence sources are at play.

In conclusion, T_1 dynamics of this qubit seems to be dominated by decoherence sources different from radioactivity, that can cause strong fluctuations both in the average value

	Source Off 1	Source On 1	Source Off 2	Source On 2
Average T_1 [μ s]	18.1 ± 0.4	19.2 ± 0.2	15.8 ± 0.3	18.6 ± 0.3

TABLE 5.3: Average T_1 values, in μ s, in the short measurements.

and in the standard deviation. Given this qubit behavior, no clear signature of the impact of radioactivity could be extracted from these datasets.

5.2.3 “Short” measurements

The fluctuating behavior of the energy-relaxation time of the qubit, with “calm” and “noisy” regimes, has been the main limit in the analysis of the data just described. A comparative study for estimating the contribution of radioactivity on that setup requires that the contribution of other decoherence sources is similar across datasets. This was not the case for our data, where qubit conditions changed multiple times across the hours.

Because of that, we decided to change our measurement strategy: instead of comparing hours of measurements, taken across multiple days, we acquired smaller sets of data, each of them lasting approximately 30 minutes. In this way, it was more likely that T_1 was stable across the datasets.

Reducing the acquisition time, though, increases the statistical uncertainty on the average T_1 , and thus reducing the discrimination power of a radioactivity-induced drop. To overcome this problem and gain a better signature, we decided to alternate measurements where the qubit was exposed to the Thorium source and measurements where the qubit was not exposed to it. In this way, if radioactivity was limiting T_1 , we would expect to observe a degradation in the qubit performance on alternate datasets. We took four datasets in total, two with the qubit exposed to the Thorium source and two without, for a total acquisition time of two hours.

The result of our measurements can be seen in Fig. 5.7. In all plots, values in red refer to datasets in which the qubit was not being exposed to the Thorium source, while values in green refer to datasets in which the qubit was exposed to it. On top, there are the superimposed histograms of the T_1 values in each dataset. To simplify comparison, three different plots have been made, comparing, respectively, the first with the second dataset, the second with the third, and the third with the fourth. On the bottom, instead, is plotted the time evolution of T_1 across the acquisition time. Average T_1 are instead reported in Table 5.3.

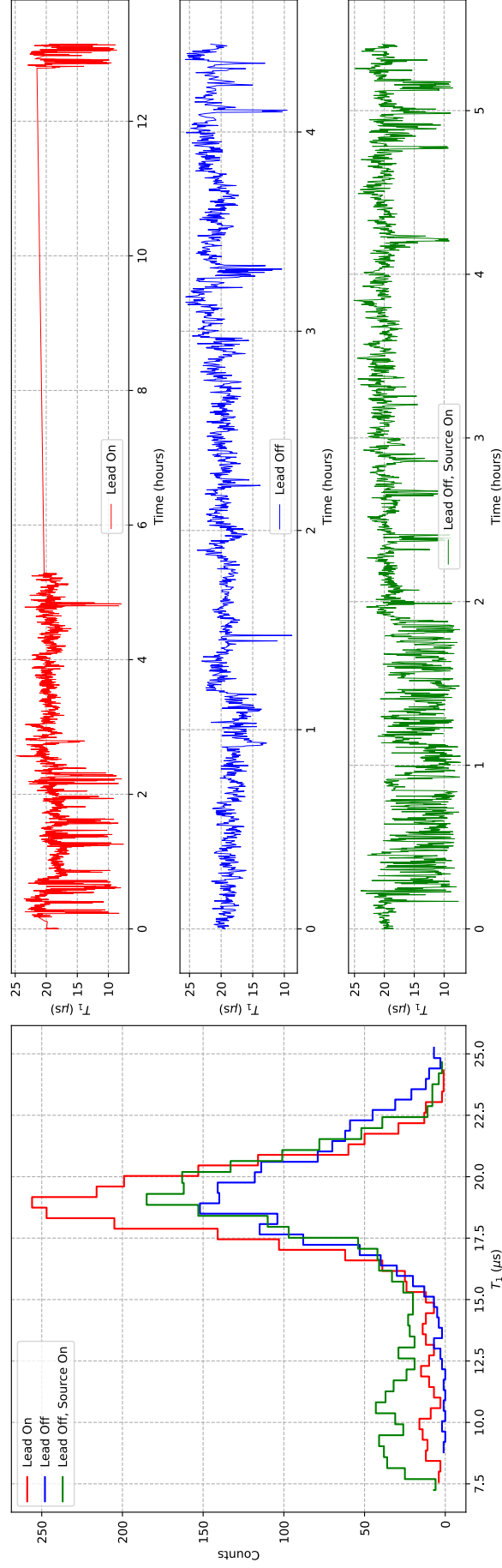


FIGURE 5.6: (Left) Histograms of the T_1 values of the fluxonium qubit measured over several hours in three different configurations: shielded against environmental radioactivity (Lead On), without the external lead and copper shield (Lead Off), and when exposed to a γ radiation source (Lead Off, Source On). (Right) The time evolution of T_1 in the three configurations. In the shielded configuration there are no data between approximately 5h and approximately 13h: this was due to an error with the software controlling the readout electronics.

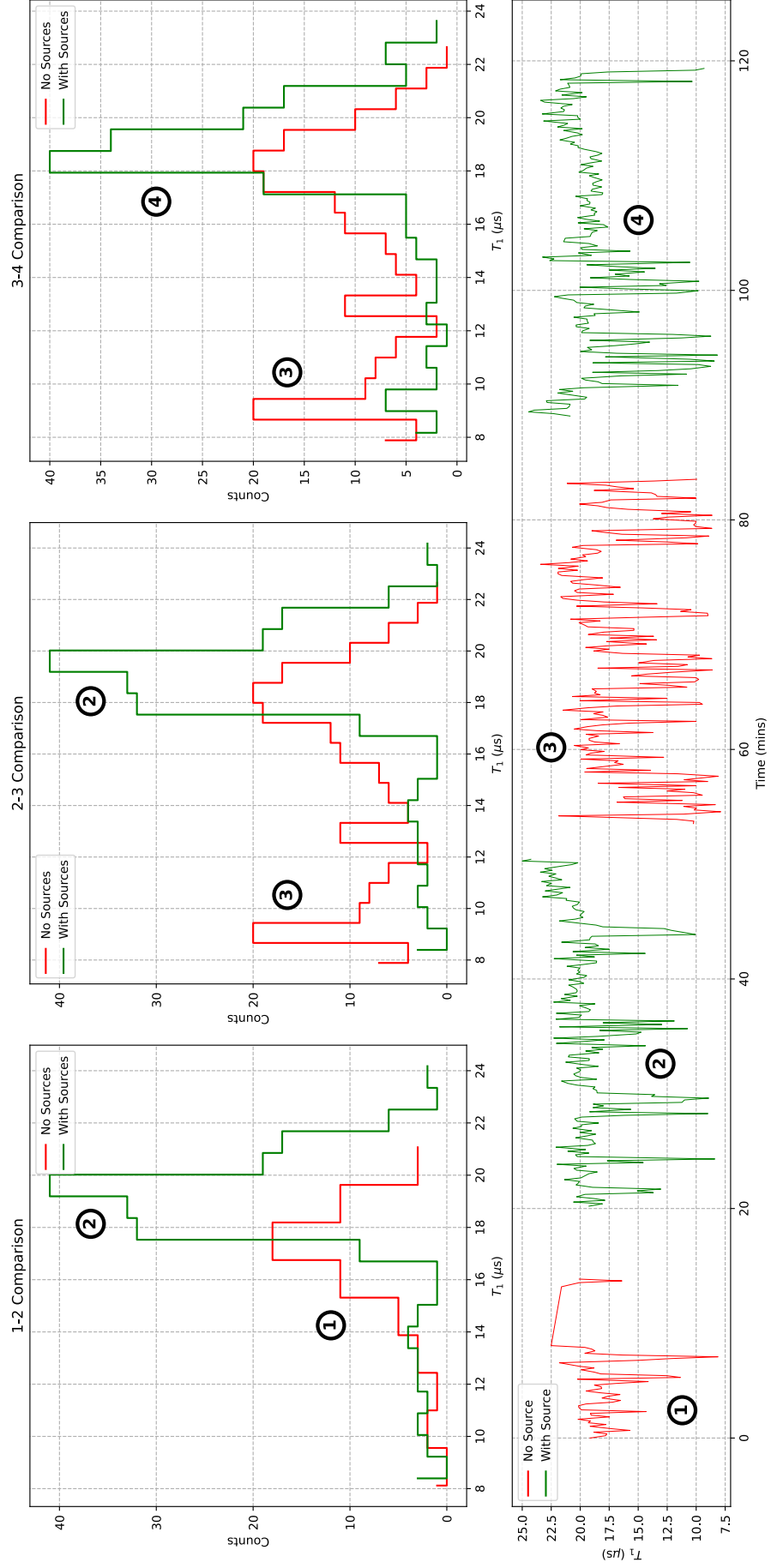


FIGURE 5.7: T_1 values of the fluxonium qubit over 2 hours, alternating periods where the qubit was exposed to a γ radiation source (red) and periods where the qubit was not exposed to it (green). (Top) Comparison of the histograms of T_1 values for consecutive datasets. (Bottom) Time evolution of T_1 .

By looking at the histograms, we observed no drop in T_1 in the measurements where the qubit was exposed to the source: the histogram featuring the biggest low- T_1 tail is actually the one of the third dataset, in which the source was not present. This is corroborated by the average values, with the datasets with the sources being the ones with the highest average T_1 . A better understanding of this finding comes from the time evolution of T_1 : during the acquisition of the third dataset the qubit entered one of its “noisy” regimes, that ended around halfway in the fourth dataset.

As for the previous measurements, this result also points to the conclusion that the dynamics of the energy-relaxation time of this fluxonium qubit is dominated by other decoherence mechanisms, with radioactivity contribution being too small to be measured with this strategy. A further confirmation of this outcome will come from the analysis of quasiparticle poisoning in the qubit, that will be described in the next section.

5.3 Looking for quasiparticle poisoning

The main advantage of measuring the energy-decay time of a qubit by looking at the rates of spontaneous excitations and decays is that, from the time distribution of the jumps, it is possible to gain further information about the decoherence mechanisms at play.

In ideal conditions, in fact, spontaneous excitations and decays are random processes, with transition probabilities depending on the overlap between the $|0\rangle$ and $|1\rangle$ eigenfunctions, on the energy difference between the two eigenstates and on the temperature of the system. Following Poisson statistics, the time separation between two consecutive jumps (or, in other words, the time the qubit spends in one of its eigenstates) should follow an exponentially decaying distribution.

It was observed [169], though, that the presence of quasiparticle poisoning alters this distribution in fluxonium qubits, causing deviations from the expected exponential. As already proved by previous experiments, described in Chapter 3, quasiparticle poisoning is actually the reason why ionizing radiation is problematic for superconducting qubits: the phonons produced by the interactions of high-energy particles with the chip substrate break the Cooper pairs in the superconductor, increasing the quasiparticle density by several orders of magnitude. When a particle deposits energy in the chip, then, we should observe a significant deviation in the jump statistics from the expected exponential distribution.

It is important to notice, however, that observing a deviation from an exponential distribution does not automatically mean that we are seeing radioactivity affecting this qubit:

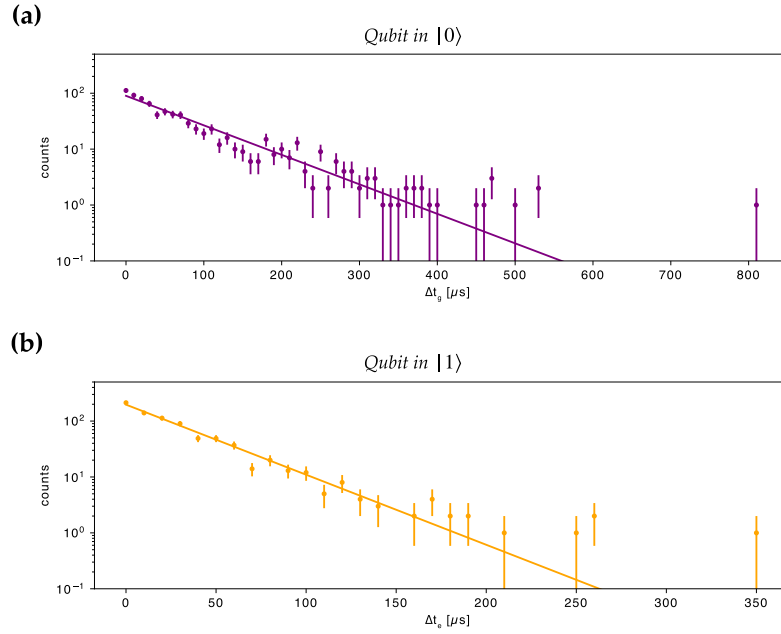


FIGURE 5.8: Histograms of the values of (a) Δt_g and (b) Δt_e for one of the traces analyzed. The line is the exponential fit. Error bars represent the statistical uncertainty estimated with Poisson statistics.

ionizing particles are not the only source of quasiparticles (infrared photons are the other major source), and there are theoretical models hypothesizing that also the coupling of the qubit with two-level systems can induce non-Poissonian behavior [170]. Also in this case, then, it is necessary to follow the comparative approach already employed in the previous analysis.

To quantify deviations from the exponential distribution we performed a χ^2 test. After dividing the data into the same traces used for T_1 measurements (so 10^6 readouts per trace), we produced for each trace two histograms: one was of the time lengths Δt_e the qubit was measured in $|e\rangle$ before jumping to $|g\rangle$, and the other of the time lengths Δt_g the qubit was measured in $|g\rangle$ before jumping to $|e\rangle$. Fig. 5.8 shows, as an example, the two histograms for one of the traces considered. Each histogram was then fitted with an exponential function and the χ^2 was computed to quantify the goodness of the fit.

The following step was conceptually identical to what was already done in the T_1 study: for each dataset we computed the distribution of the χ^2 values and their time evolution, in order to compare datasets in which the qubit was exposed to ionizing radiation and datasets in which it was shielded. If radioactivity was increasing significantly the quasiparticle poisoning of the qubit, this would have resulted in a stronger deviation of the Δt_g and Δt_e histograms from the exponential distribution, and, hence, in a higher χ^2 value.

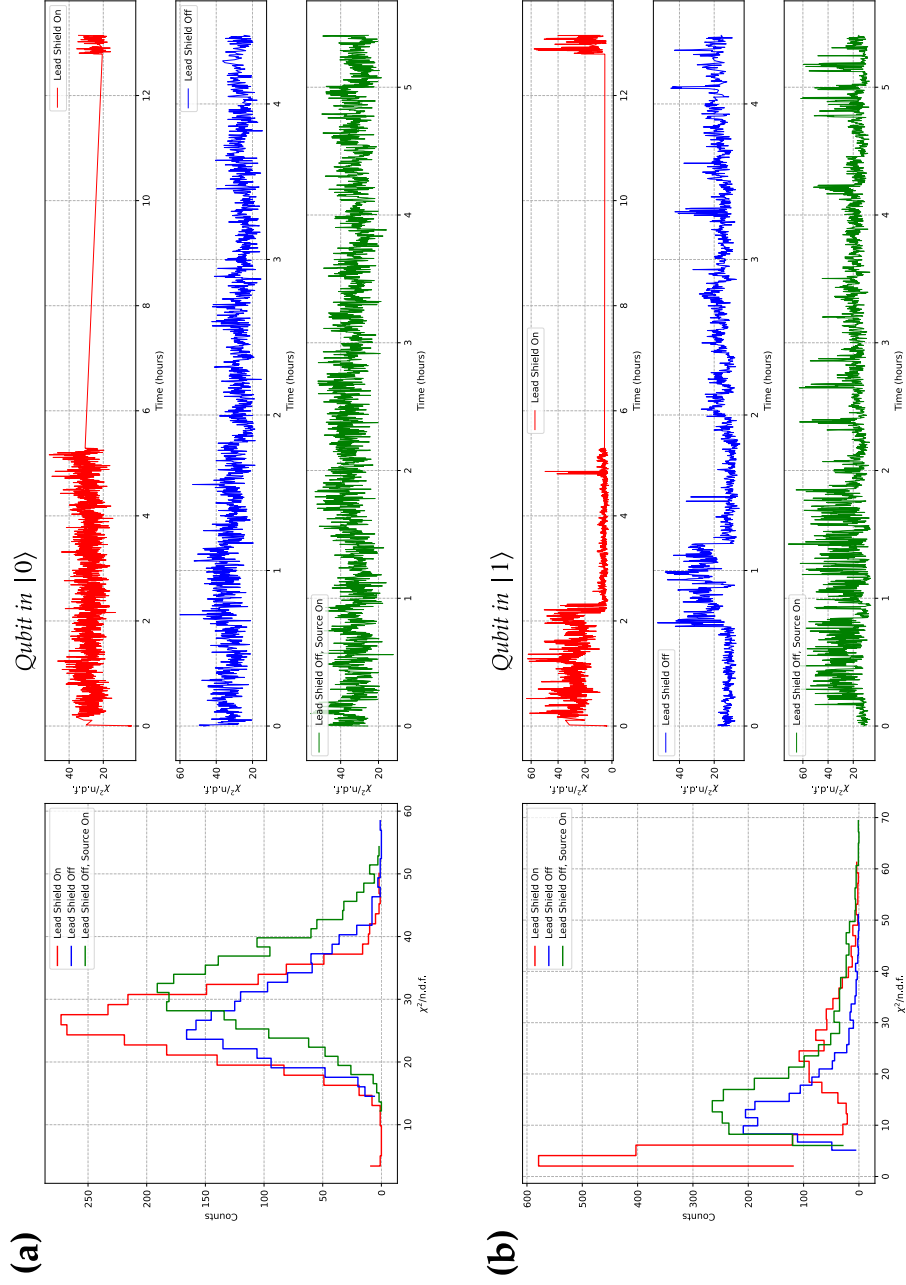


FIGURE 5.9: χ^2 values obtained when fitting the distribution of time intervals the qubit spent (a) in the ground state and (b) in the excited state, before jumping to the other state. On the left, histograms of the χ^2 values for the three “long” datasets: qubit fully shielded against radioactivity (in red), qubit without the external lead shield (in blue) and qubit exposed to a Thorium radioactive source (in green). On the right, time evolution of the χ^2 for the same datasets.

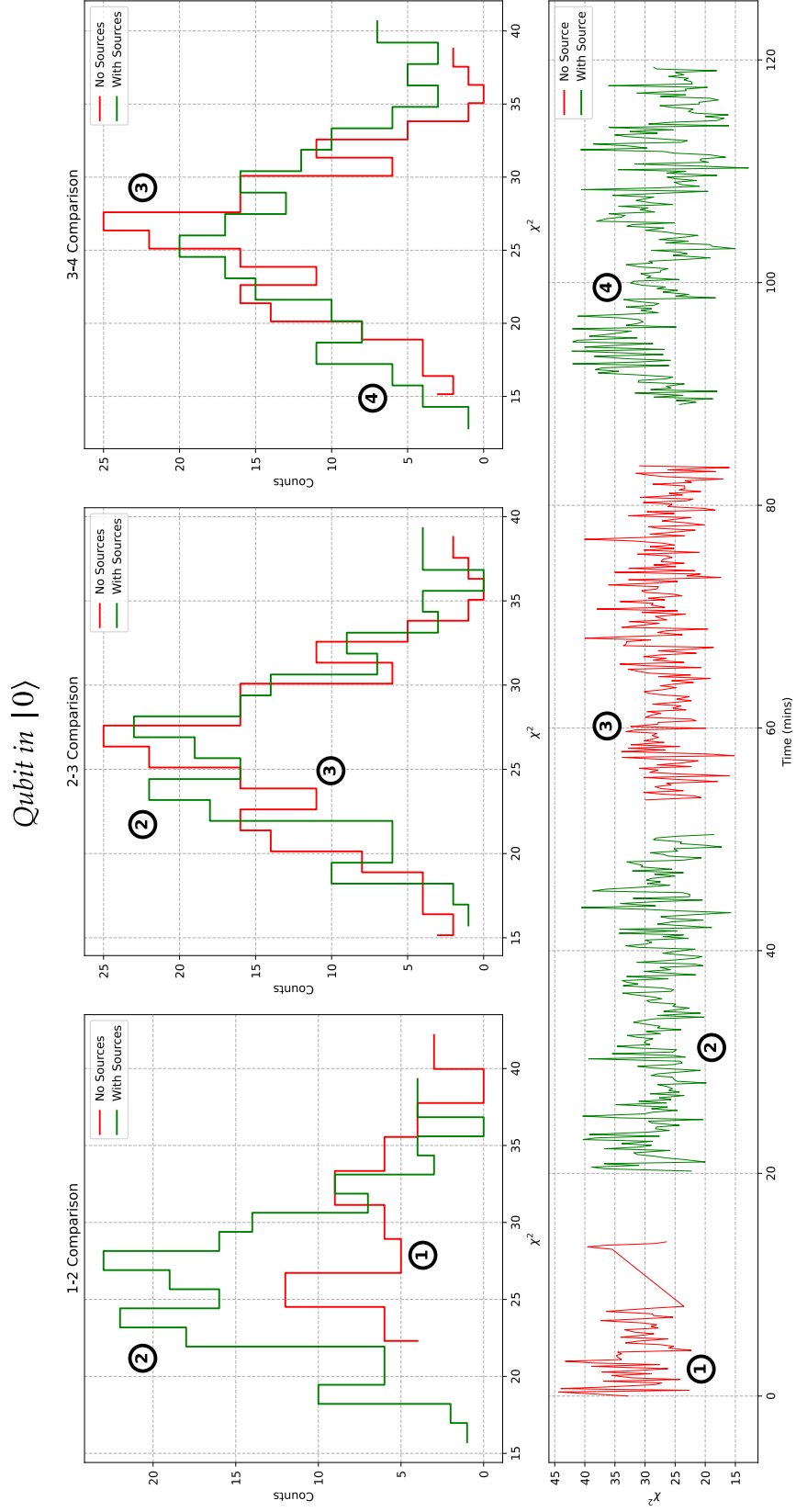
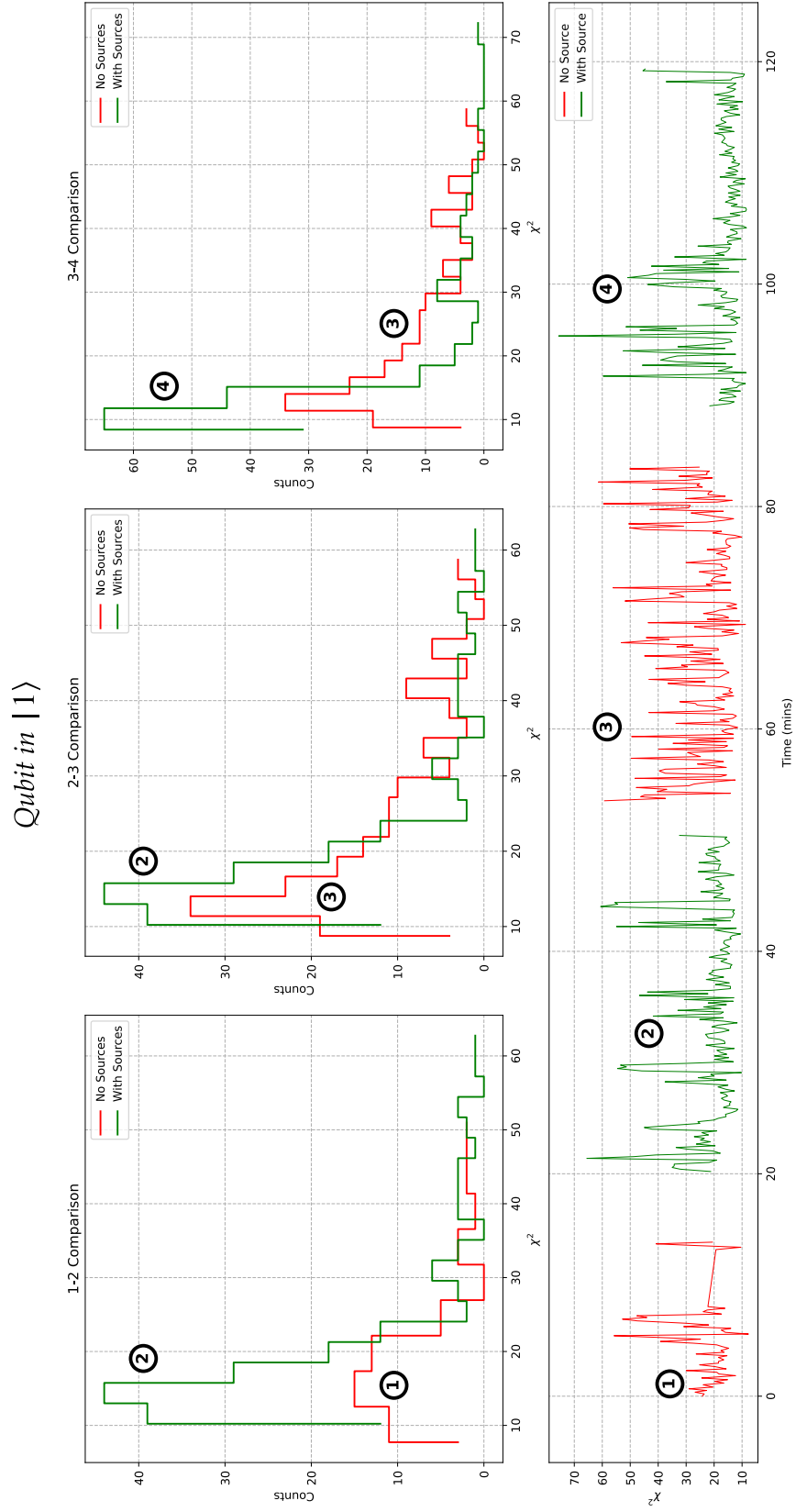


FIGURE 5.10: χ^2 values for the fit of the Δt_g distribution over 2 hours, alternating periods where the qubit was not exposed to a γ radiation source (red) and periods where the qubit was exposed to it (green). (Top) Comparison of the histograms of the χ^2 values for consecutive datasets. (Bottom) Time evolution of χ^2 .

FIGURE 5.11: Same as Fig. 5.10, but for the fit of the distribution of Δt_e .

	Average $\chi^2/\text{n.d.f.}$		
	Lead On	Lead Off	Lead Off, Source On
Qubit in $ 0\rangle$	28.1 ± 0.1	29.3 ± 0.2	33.6 ± 0.2
Qubit in $ 1\rangle$	15.3 ± 0.3	16.5 ± 0.2	20.9 ± 0.2

TABLE 5.4: Average values of the reduced χ^2 in the long measurements

	Average $\chi^2/\text{n.d.f.}$			
	Source Off 1	Source On 1	Source Off 2	Source On 2
Qubit in $ 0\rangle$	31.8 ± 0.7	27.8 ± 0.4	26.8 ± 0.3	27.8 ± 0.5
Qubit in $ 1\rangle$	24 ± 2	22.0 ± 0.8	25.0 ± 0.9	10.8 ± 0.8

TABLE 5.5: Average values of the reduced χ^2 in the short measurements

This analysis was performed on the same datasets used for the study of T_1 , both the “long” ones, lasting a few hours, and the “short” ones, lasting around 30 minutes. Concerning “long” measurements, the results are summarized in Fig. 5.9, while the average values of the reduced χ^2 for each dataset are reported in Table 5.4. Concerning Δt_g (Fig. 5.9a), the average value of the reduced χ^2 slightly increases once the external lead shield is removed, and increases further when the qubit is exposed to the Thorium source. The same effect seems to be much more evident when looking at the average values of the reduced χ^2 for the fit of the Δt_e distributions. If, however, we look at the time evolution in the three datasets we can see that, as for the T_1 study, the qubit alternates “calm” periods with “noisy” ones, again undermining the hypothesis that radioactivity could be the cause of that discrepancy.

The same analysis was repeated by looking at the “short” datasets, both for Δt_g and for Δt_e . The results of the analysis are pictured in Fig. 5.10 and Fig. 5.11, with average values reported in Table 5.5. Unfortunately, the variation of the average value of the χ^2 that we observed when comparing the shielded qubit with the qubit exposed to the Thorium source was not so evident in short measurements. If we look at Δt_g , we can see that we have the highest χ^2 values in the first short dataset, when the qubit is actually shielded against radioactivity. In the next three datasets, instead, the histograms of the χ^2 values with and without the Thorium source mostly overlap, without showing the shift that we observed, instead, in the long measurements. We got a similar outcome also when looking at Δt_e , with the highest χ^2 values reached in the third dataset, again one in which the qubit was shielded.

Also for this analysis, then, we concluded that the main source of non-Poissonian quantum jumps is not radioactivity, and that other phenomena are actually dominating the dynamics of the qubit under study.

5.4 Summary of the results

Summarizing what has been written in the past sections, we concluded that ionizing radiation does not significantly affect the decay time of the fluxonium qubit we tested. Exposing the qubit to ionizing radiation did not result in a statistically significant worsening of its average T_1 , compared to what was obtained when the qubit was fully shielded. At the same time, deviations of the quantum jumps distributions from the expected Poissonian behavior, which are known to be correlated also with quasiparticle poisoning, do not vary significantly between datasets, giving similar results when the qubit is shielded and when it is exposed to the radiation source.

Both these results lead to the conclusion that qubit decoherence is dominated by phenomena different from radioactivity. However, this does not necessarily mean that radiation is not affecting the qubit at all: its contribution may just be marginal compared to the other decoherence sources.

A possible hint of why no significant effect of radioactivity was measured comes from Wilen et al. [131]: going back to Fig. 3.3c, it is possible to see that the fluctuations in the decay time of their qubits that they associated to particle impacts lasted only on a millisecond scale, with the qubit going back to its usual behavior after it. On the contrary, a T_1 measurement like the one we performed in our experiment lasted a few seconds, because in order to estimate correctly $\Gamma_\uparrow$ and $\Gamma_\downarrow$ we needed to acquire a high statistics of quantum jumps. Such a rapid fluctuation, when counting over such a long timescale, may just not affect the final result of the measurement.

In order to prove this hypothesis true, though, it is necessary to change the measurement strategy, employing an algorithm for the measurement of T_1 fast enough to be sensitive to fluctuations at the millisecond timescale. This was the aim of our next experiment, performed on a transmon qubit as part of the activities of the Superconducting Quantum Materials and Systems (SQMS) Center, and will be described in the next chapter.

Chapter 6

Detecting gamma radiation with transmon qubits

The experiment with the fluxonium qubit described in Chapter 5 showed that, for qubits with decay times of tens of μs , ionizing radiation is not a significant limiting factor. Despite this result, there are several reasons why the impact of ionizing radiation should be addressed in the development of quantum computers. The most obvious one is that, as the design and production of superconducting qubits improve, and other sources of decoherence are mitigated, environmental radioactivity remains, eventually setting an upper limit on the achievable decay time. The simulation framework described in Chapter 4 showed that, for quantum chips with $O(\text{cm}^2)$ surface, it is reasonable to expect a particle impact every few seconds, meaning that any computation on that chip requiring a longer time is expected to fail even if other decoherence sources are mitigated. For larger chips, which are necessary as quantum computers increase in computational power, impacts will be even more frequent.

The second reason is that ionizing radiation can cause quantum error correction algorithms, which would be the computational solution for the limit above described, to fail. As mentioned in Section 2.3 and then discussed extensively in Section 3.2, particle impacts produce correlated errors that span across multiple qubits, with quasiparticle densities orders of magnitude higher than what would be bearable by error correction algorithms, thereby making this approach difficult to pursue with present-day technologies.

To address these concerns, though, it was necessary to change our measurement strategy: a T_1 measurement like the one described in Chapter 5, in fact, proved to be unsuitable for this study, because of the long time required – around three seconds – to obtain a reliable estimate.

The reason why that T_1 measurement was too slow for detecting particle impacts in the chip can be hinted from the analysis done by Wilen et al. [131] and by McEwen et al. [133]. In their studies, quasiparticle bursts associated with radiation events produced drops in the decay time lasting only a few milliseconds. This can be explained by the fact that, in every superconductor, quasiparticles produced in these transient events would recombine to form Cooper pairs again or be trapped into magnetic vortices, until the equilibrium population is restored.

Considering that, even when the qubit was exposed to the Thorium source, we would expect less than ten impacts per trace, the result was that the qubit was perturbed for a time three orders of magnitude shorter than the one required for a T_1 measurement, while behaving normally the rest of the time. The T_1 value we measured would then be insensitive to such a fast fluctuation, resulting in the outcome described in the last chapter.

In order to verify this hypothesis and directly detect radiation-induced decays, we conducted a new experiment. This time, we modified the measurement scheme already employed by McEwen et al. [133], adapting it to a single qubit and changing the reset protocol to a different one that can work with fixed-frequency qubits.

This experiment was conducted as part of the activities of the SQMS Center on a transmon qubit developed by the collaboration. The qubit was measured in two different locations: underground at INFN-LNGS and, above-ground, at Fermilab, where the SQMS headquarters are located.

At LNGS we again employed the Thorium sources used in the previous experiment, in order to compare the behavior of the qubit in a shielded configuration and when exposed to radioactivity. This time, though, we decided to employ sources with different activities, with the idea that, if the detected events are actually radiation-induced, we should see their rate increase with the activity of the source. To simplify operations, the sources were placed between the dilution refrigerator and the external shielding, so that it was not necessary to remove and reinstall the lead shield.

The above-ground measurements at FNAL, instead, allowed us to compare underground measurements with those obtained in a typical laboratory environment, with the qubit unprotected both from environmental γ -rays and from cosmic muons.

The results of this experiment, along with a detailed description of the measurement and analysis process, have been summarized in a recent article [171] titled *Evaluating radiation impact on transmon qubits in above and underground facilities*, published on *EPJ Quantum Technology*.

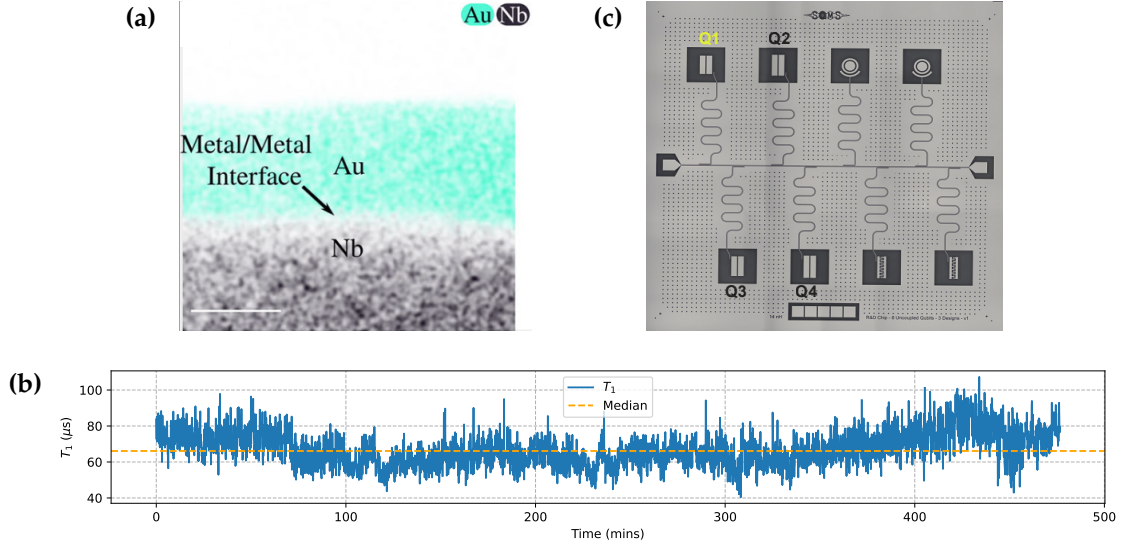


FIGURE 6.1: (a) Section, in false colors, of the gold capping encapsulating the Niobium superconducting island of the qubit, taken from [90]. (b) T_1 values of the qubit used in this study measured across approximately 8 hours. (c) Photo of the chip employed for this experiment; the qubit chosen is labeled Q1; other qubits were discarded due to their worse performance in terms of T_1 and readout fidelity.

6.1 Experimental setup

6.1.1 The qubit

The qubit that we used is a fixed-frequency transmon qubit developed by the SQMS Center [89], which uses niobium as superconducting material. The advantage of using niobium instead of aluminum lies in its higher critical temperature, which results in a higher superconducting gap. The higher gap is convenient for reducing quasiparticle-induced decoherence, but, as will be proved in this chapter, is not sufficient to make the qubit insensitive to the high-energy phonons produced by a particle impacting the chip.

One of the main sources of decoherence in Niobium qubits is the oxidation of their surfaces when exposed to air. This is because Niobium oxides, being amorphous, contain a high density of two-level systems that can interfere with qubit operations. To prevent the formation of oxides, the surface of the qubit is encapsulated with a thin layer of gold evaporated on it approximately 10 nm thick (see Fig. 6.1a). The suppression of this loss mechanism provided by the surface encapsulation allowed to reach, for the qubit under study, T_1 values between 50 and 120 μs (see Fig. 6.1b), but qubits with T_1 values ranging up to 600 μs have also been fabricated [89]. The frequency of the readout resonator was 7.207 GHz, while the transition frequency ω_{01} was 4.722 GHz.

The qubit is fabricated on a $7.5 \times 7.5 \text{ mm}^2$ chip with $437 \mu\text{m}$ thickness (see Fig. 6.1c), along with seven other qubits that were not included in this study due to their lower T_1

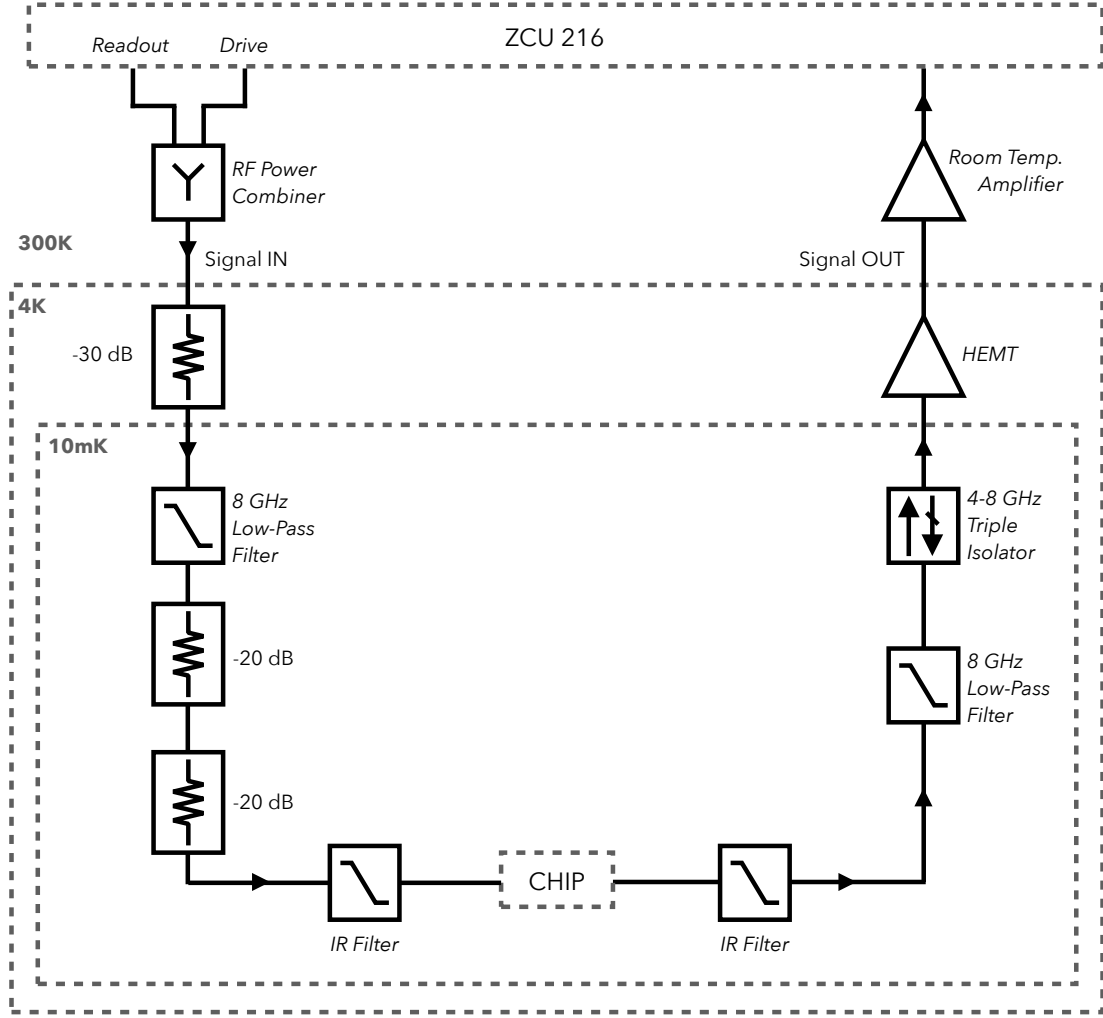


FIGURE 6.2: Scheme of the RF lines and electronics employed for the experiment. Here, readout and drive signal reach the qubit through the same RF line, so the two outputs of the ZCU216 are combined at room temperature before entering the cryostat.

and readout fidelity. The chip consists of a sapphire substrate covered in niobium, that acts as a ground plane, traversed by a transmission line used both for readout and drive and with the eight qubits with their readout resonators capacitively coupled to it. Two ports, at the ends of the transmission line, are used for input and output respectively.

6.1.2 RF lines and room-temperature electronics

The scheme of the RF lines and components employed in this experiment is portrayed in Fig. 6.2. The idea is the same as the one employed for the experiment with the fluxonium: the qubit was placed at the coldest stage of the dilution refrigerator, where a base temperature of approximately 10 mK is reached.

Readout and drive signals were produced at room temperature and then attenuated and filtered along the input line. Output readout signals were amplified in two stages: the first one is the HEMT amplifier at 4K, while the second one is the room-temperature amplifier already employed in the experiment with the fluxonium. The overall amplification is approximately 75 dB, to be compared with a total attenuation of 70 dB of the input signal (plus approximately 15 dB more of attenuation at 7 GHz due to the lines). No parametric amplifier at 10 mK was used in this experiment. That is because, in the beginning, a Josephson Parametric Amplifier was also initially planned. Unfortunately, though, the DC line used for its biasing introduced extra noise disturbing the qubit measurement, so we decided to remove it. Nonetheless, we were able to discriminate qubit states with readout fidelity above 90%, which was sufficient for our purpose.

For the generation and the demodulation of RF signals we employed a Xilinx ZCU216 RFSoc board. The board is equipped with the Quantum Instrumentation Control Kit (QICK) firmware [172] for the generation and the timing of the pulses and controlled via a Python interface. This choice was motivated by the flexibility offered by the board and by its firmware, that allowed us to implement the measurement protocol that will be described in Section 6.2.

6.1.3 Shielding and expected contributions from radioactivity

The shielding employed for the low radioactivity measurements is the same as that used in the fluxonium experiment. The qubit, inside its copper box, is first shielded by a CryoPermTM magnetic shield and is located below the mixing chamber. Directly above it, a 3-cm-thick lead shield protects it against γ radiation coming from above (see Fig. 6.3a).

Outside the dilution refrigerator, a second magnetic shield, made of two Permimphy layers, was placed, surrounded by the lead and copper shield (see Fig. 6.3b). The copper shield was made of copper bars with 2 cm thickness and $10 \times 66.5 \text{ cm}^2$ surface, while for the lead shielding we used two rounds of $10 \times 20 \times 5 \text{ cm}^3$ bricks.

In order to better estimate the rate of particle impacts in the chip, we also improved the Geant4-based simulations already described in Chapter 4, replacing the single cylindrical lead shield with a more realistic geometry made of bricks and replacing the simplified chip model, including its copper box, with a realistic representation of the configuration used in this experiment. The updated geometry can be seen in Fig. 6.3c.

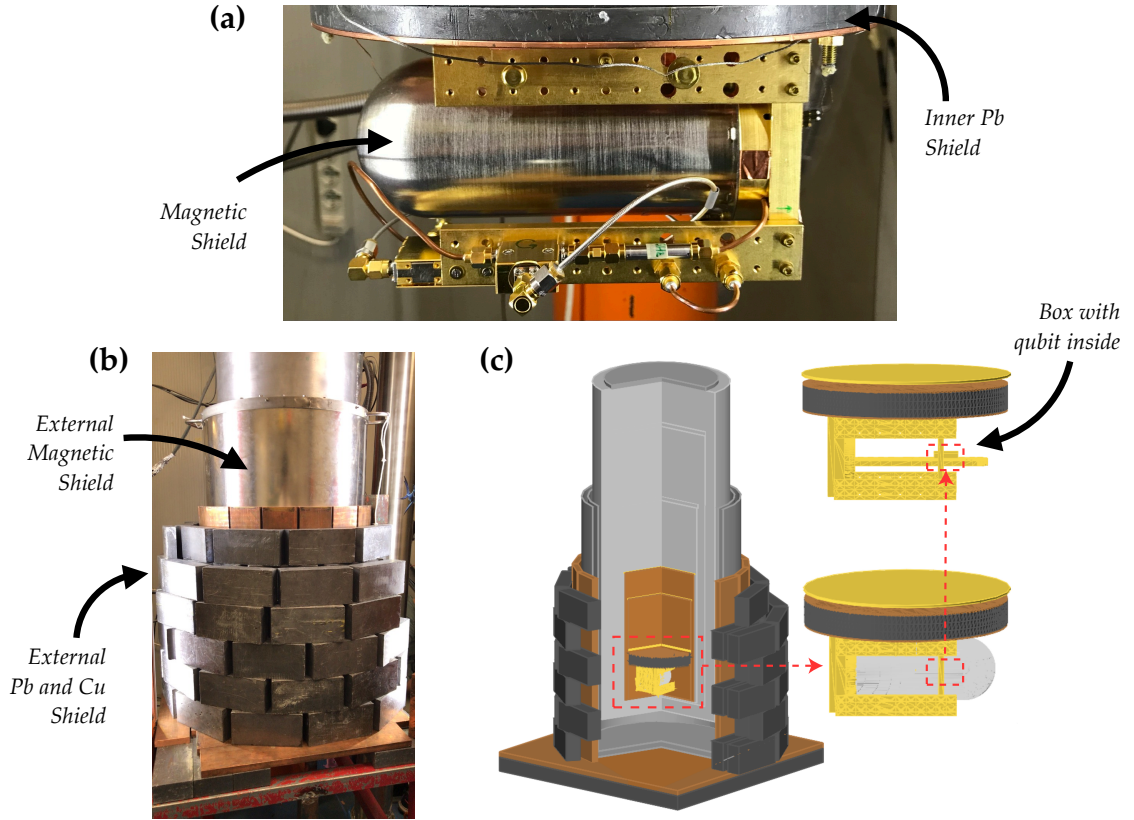


FIGURE 6.3: (a) Photo of the magnetic shield, containing the box with the qubit inside, with the 3-cm-thick inner lead shield on top of it. (b) Photo of the external shielding, with the second magnetic shield and the external lead and copper shielding. (c) Reconstruction of the experimental setup for the Geant4-based simulation.

With this update, we estimated contributions from the most impacting sources found in Chapter 4, so environmental γ 's, cosmic muons and two Printed Circuit Boards, located at opposite ends of the chip.

The estimation of impact rates followed the same procedure already described in Chapter 4, starting from the measured fluxes reported in Section 4.2.1.

From our simulations, we estimated that the overall rate at LNGS, when the lead shielding is implemented, is one order of magnitude lower than in FNAL, 4×10^{-3} events/sec against 42×10^{-3} events/sec. When employing Thorium sources, instead, we estimated the rate at LNGS to increase up to two orders of magnitude, with 0.43 ± 0.05 events/sec in the case of the 161 kBq source, the most active one. Results are summarized in Tables 6.1 and 6.2. Fig. 6.4, instead, shows the spectra of the energy deposited in the chip in the two locations (Fig. 6.4a) and the spectra of the energy deposited by the Thorium sources employed at LNGS (Fig. 6.4b).

Source	LNGS [events/sec]	FNAL [events/sec]
Env. γ -rays	$(1.3 \pm 0.1) \times 10^{-3}$	$(31 \pm 2) \times 10^{-3}$
Muons	$< 10^{-5}$	$(8 \pm 0.5) \times 10^{-3}$
Setup	$(2.7 \pm 0.5) \times 10^{-3}$	$(2.7 \pm 0.5) \times 10^{-3}$
Total	$(4.0 \pm 0.6) \times 10^{-3}$	$(42 \pm 3) \times 10^{-3}$

TABLE 6.1: Estimated contributions, in terms of impacts per second, of environmental radiation sources in the two experimental locations, obtained from simulations.

Thorium source activity [kBq]	Rate [events/sec]
44.2	0.12 ± 0.01
75.9	0.20 ± 0.02
125.4	0.34 ± 0.04
161.0	0.43 ± 0.05

TABLE 6.2: Expected contributions, in terms of impacts per second, of the Thorium sources employed in the experiment.

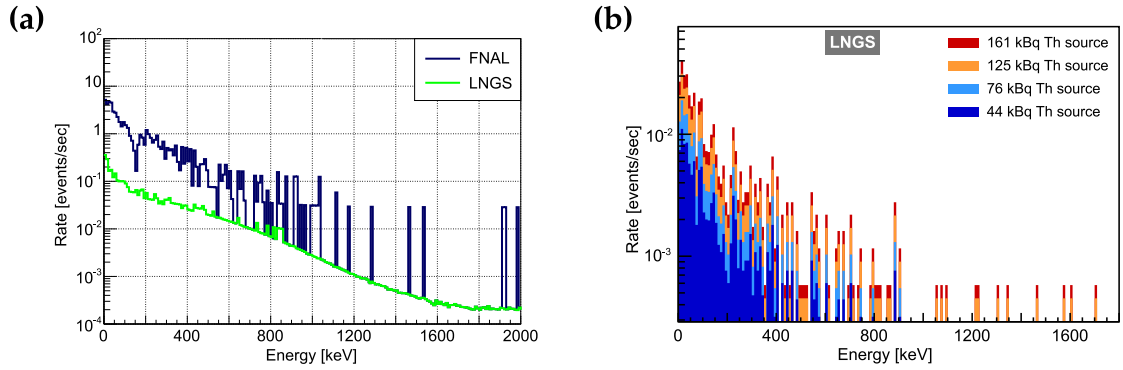


FIGURE 6.4: (a) Spectrum of energy deposited in the chip, according to the simulations, at LNGS (green) and at FNAL (blue). (b) Spectrum of the energy deposited in the chip by the γ 's produced by the Thorium sources, estimated also with simulations.

6.2 Measurement strategy

As already mentioned in the introduction to this chapter, a standard T_1 measurement could be too slow to detect the quasiparticle poisoning produced by ionizing radiation, that Wilen et al. [131] and McEwen et al. [133] measured to last a few milliseconds at most. For our experiment, it was therefore necessary to employ a measurement strategy sensitive to these short-lived fluctuations in the decay time.

In their experiment, both Wilen et al. and McEwen et al. employed a similar detection scheme: the qubits were first excited to $|1\rangle$ and, then, measured after a delay of a few microseconds. Fluctuations in the decay time of the qubits would result in a varying

probability of measuring the qubit still in $|1\rangle$ after the delay. In the absence of quasiparticle poisoning, that probability would be high, whereas it would drop to very low values when the qubit was disturbed by particle interactions.

Directly implementing one of these schemes in our experiment, however, would have been problematic for two reasons, both concerning state preparation. In Wilen et al., in fact, a π -pulse was sent to the qubit in order to prepare it in $|1\rangle$. Yet, as seen in Section 2.1, if the qubit is already in $|1\rangle$, applying a new π -pulse would send its state to $|0\rangle$. These preparation errors significantly complicate a lot the estimation of the decay time and, hence, the detection of particle impacts, while waiting longer between measurements, so as to be almost certain that the qubit had decayed in the meantime, would have drastically reduced the number of measurements that we could do during the quasiparticle poisoning. McEwen et al. solved this issue by improving the detection scheme, introducing a reset protocol that first set the qubit state to $|0\rangle$, regardless of its original state, and only after that sends the π -pulse. Nonetheless, that reset protocol relies on rapidly changing the qubit transition frequency, something that could not be done on our fixed-frequency qubit.

In order to overcome these issues, we further modified the measurement scheme introduced in these previous works by implementing a reset protocol based on a *conditional* π -pulse. The qubit state was first measured to set it in one of its eigenstates. Immediately afterwards, a π -pulse was sent only if the measured state was $|0\rangle$, preventing unwanted $|1\rangle \rightarrow |0\rangle$ transitions.

The complete measurement scheme is shown in Fig. 6.5a. The qubit was first prepared in $|1\rangle$ through the reset protocol just described. Then, its state was measured again after a few microseconds. In normal conditions, the waiting period between state preparation and measurement would be much smaller than the T_1 of the qubit. By repeating this measurement, then, we would observe in the data a continuous stream of 1 outcomes. In the presence of quasiparticle poisoning, instead, the T_1 would drop to values shorter than the waiting period, resulting in a high probability of measuring 0. In this scenario, the data stream exhibit a sudden change, with a stream of 0's lasting until the quasiparticle density returns to its equilibrium value. Fig. 6.5b depicts an event observed in our data, in which such a strong transient fluctuation of the decay time is clearly visible: for the first approximately 7 milliseconds the qubit is behaving normally, and we measured mostly 1; suddenly after that, its response changes dramatically, and for all the subsequent measurement the measured state is 0. The qubit, then, returns to its normal behavior within a couple of milliseconds.

For our experiment, the waiting period was chosen to be equal to $5\ \mu\text{s}$. With this choice, the probability $P(0)$ of observing the qubit in $|0\rangle$ after the state preparation and the

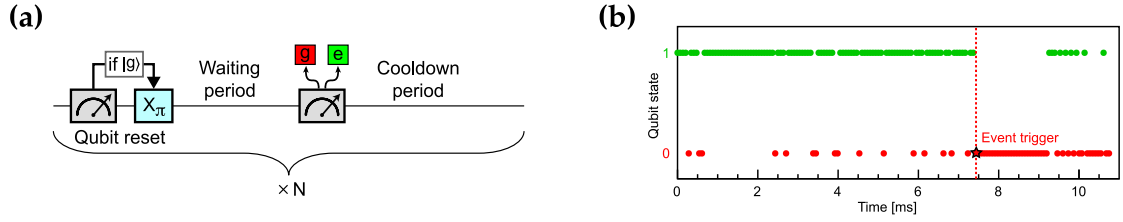


FIGURE 6.5: (a) The measurement scheme employed in our experiment: a reset protocol based on a conditional π -pulse prepares the qubit in $|1\rangle$; the qubit state is then measured after a waiting period of a few μs (for our qubit we found $5 \mu\text{s}$ to be the best value); a cooldown period of a few tens of μs after the second readout prevents state leakage. (b) Example of a detected event compatible with radiation-induced quasiparticle poisoning: in normal conditions, between 0 and $7.5 \mu\text{s}$, the qubit is behaving normally and its state is measured to be $|1\rangle$ most of the times; at $7.5 \mu\text{s}$, however, the decay time of the qubit drops, resulting in a stream of $|0\rangle$'s in the data lasting a few microseconds.

delay has the largest discrepancy between the normal scenario and the scenario in which the qubit is disturbed by radioactivity. By calling T_1 the typical relaxation time of the qubit and T_1^{qp} the relaxation time in the presence of quasiparticle poisoning, the discrepancy of $P(0)$ is given by:

$$\Delta[P(0)] = \exp(-t_d/T_1) - \exp(-t_d/T_1^{qp}) \quad (6.1)$$

Where t_d is the waiting period. This difference is maximum when:

$$t_d = \frac{T_1^{qp} \cdot T_1}{T_1 - T_1^{qp}} \ln \left(\frac{T_1}{T_1^{qp}} \right) \quad (6.2)$$

In Section 6.1.1 it was shown that T_1 of this qubit exhibits significant temporal variations, with a median value slightly below $70 \mu\text{s}$. At the same time, results from McEwen et al. showed that, for their qubits, T_1^{qp} was lower than $1 \mu\text{s}$. Using $70 \mu\text{s}$ and $1 \mu\text{s}$ as rough estimates for T_1 and T_1^{qp} , respectively, the optimal waiting period lies between 4 and $5 \mu\text{s}$, that we approximated with $5 \mu\text{s}$.

After the state readout, we included a *cooldown period* of a few tens of microseconds in the measurement cycle, in which no operation was done on the qubit (see Fig. 6.5a). This was necessary because, by continuously sending pulses to the qubit, we inadvertently heated it, inducing transitions also to the second and the third excited state. Adding that extra waiting time allowed us to minimize the state leakage probability. The cooldown period, however, could not be increased indefinitely: a longer delay results in fewer measurement cycles in a millisecond window, thus reducing our capability to identify radiation-induced events. A systematic study of how the duration of the cooldown period affects the experiment has been performed and will be described in Section 6.4.2.

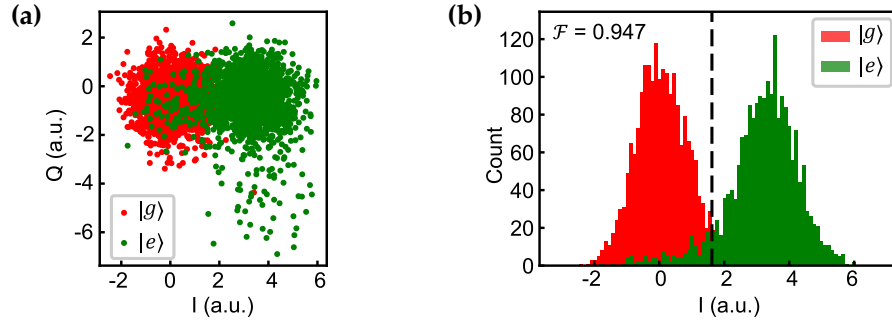


FIGURE 6.6: Example of a fidelity measurement in our experiment. (a) The I – Q clouds associated with $|0\rangle$ and $|1\rangle$, here labelled as $|g\rangle$ and $|e\rangle$, respectively. (b) Histogram of the I values for the two clouds. The fidelity reported is computed as $\mathcal{F} = 1 - \frac{1}{2}[P(0|1) + P(1|0)]$

In addition to the waiting period and the cooldown period described above, three other time lengths determined the overall sampling time of our experiment: the time durations of the two readouts and of the π -pulse, with the former varying between 5 and 9 μs and the latter of approximately 200 ns, and the *ringdown time* of the readout resonator. The ringdown time is the characteristic decay time of light inside a resonant cavity: it is particularly important when implementing the qubit reset because the conditional π -pulse is sent right after the readout. If that happens before the readout pulse has fully decayed in the resonator, the two pulses will interfere, preventing proper qubit excitation. For our chip, that meant waiting extra 500 ns between the first readout and the subsequent conditional π -pulse.

Overall, we acquired data with sampling times T_S varying between 39.8 μs and 73.6 μs . In particular, datasets with T_S of 67.6 μs and 73.6 μs were acquired at both LNGS and FNAL, while datasets with other sampling times were acquired in FNAL for the systematic study on how T_S affects the detection efficiency.

6.3 Analysis of the data

6.3.1 State discrimination and leakage

To simplify state discrimination, the phase of the readout pulse was chosen so that the I – Q clouds corresponding to $|0\rangle$ and $|1\rangle$ were aligned along the Q axis, as shown in Fig. 6.6. In this way, as already done for the fluxonium experiment (Section 5.2.1), a threshold value on I is sufficient to discriminate between the two states.

The two clouds, however, are not completely separated, so chosen threshold value results in a finite classification error (i.e. a readout fidelity lower than 100%). The threshold value chosen was the one minimizing this error, so maximizing the readout fidelity. To

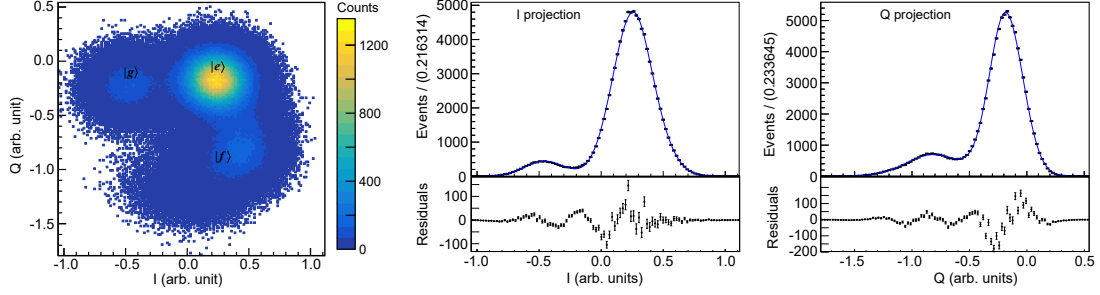


FIGURE 6.7: (Left) 2D histogram of the $I - Q$ values of the readout measurement for a single trace (10^6 measurement cycles). The states $|0\rangle$, $|1\rangle$ and $|2\rangle$, there labeled as $|g\rangle$, $|e\rangle$ and $|f\rangle$, respectively, are clearly visible. In most of the measurements the qubit is found in $|1\rangle$ due to the reset protocol. (Center and Right) Simultaneous fit of the projections along the two axes for estimating state population.

this aim, the two clouds were first fitted with Gaussian functions. Then, we varied the threshold value over the interval defined by the two Gaussian centers. For each value, we integrated the tails of the two Gaussians falling on the wrong side of the threshold, thereby estimating $P(0|1)$ and $P(1|0)$. We then chose the threshold value that minimized $P(0|1) + P(1|0)$. In this way, the typical fidelity values achieved were of 90-95%.

The continuous probing of the qubit state, though, introduced another complication: the leakage out of the computational space of the qubit state. Although the aim of the cooldown period in our detection protocol was to prevent unwanted transitions to excited states higher than $|1\rangle$, state leakage could not be completely avoided. This poses a challenge for data consistency, because characteristic decay times vary if the qubit is in $|2\rangle$ or $|3\rangle$ instead of $|1\rangle$. We therefore excluded data traces (each trace corresponds to 10^6 measurement cycles) with high $|2\rangle$ and $|3\rangle$ populations.

This approach guaranteed us datasets with similar populations of the qubit states, which, as we found out, is a fundamental requirement for our analysis. The impact of fluctuating populations across datasets will be investigated more in detail in Section 6.4. Satisfying this requirement, however, comes at a cost: it reduced the effective live-time of our measurements, as can be seen in Table 6.3, that summarizes all the datasets acquired.

To identify and reject traces with high leakage, we produced, for each of them, a 2D histogram of the $I - Q$ values of the readout pulses. The two projections of this histogram along the quadratures were then simultaneously fitted using the RooFit toolkit [173]. For each projection, the fit function was a sum $\sum_i A_i \cdot f(x, \bar{x}_i, \sigma_i)$ of Gaussian functions centered at $\bar{x}_i$ with standard deviation σ_i , with x corresponding to I or Q . A_i are the normalization factors, representing the number of events per each state i , and were shared between the two fit functions. For each state i , its population was then estimated as $P_i = A_i / \sum_i A_i$.

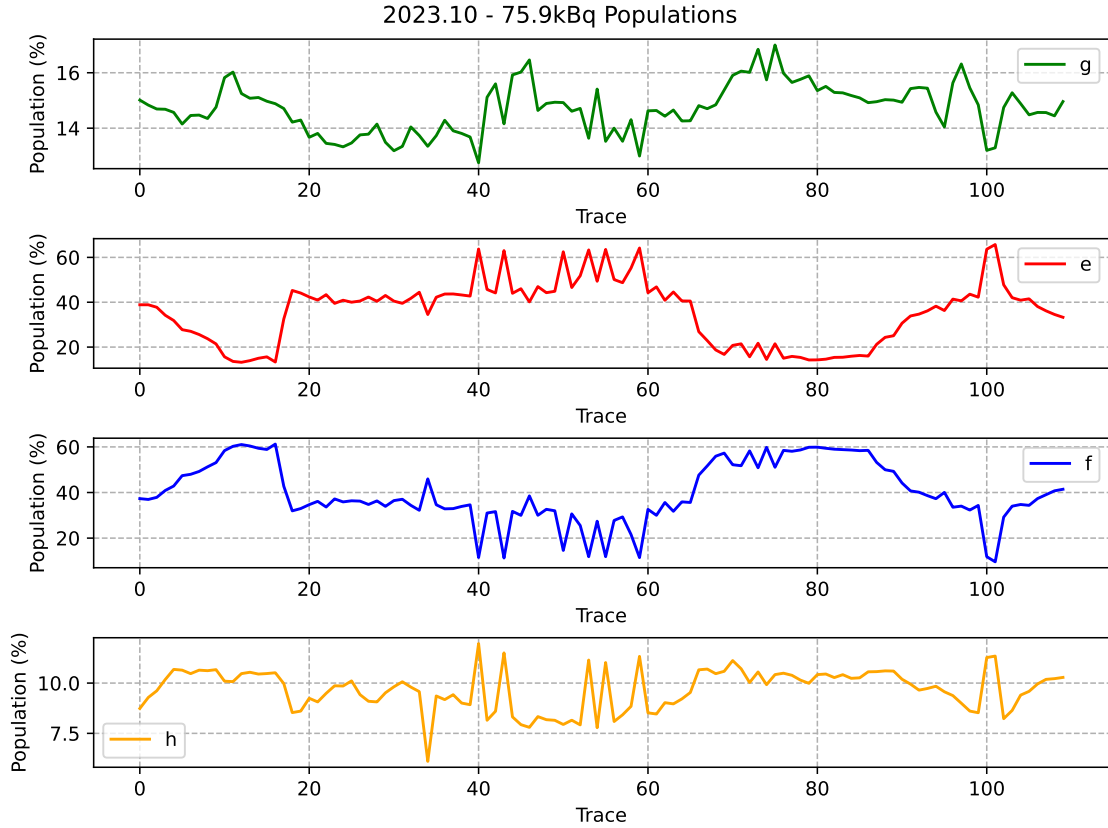


FIGURE 6.8: State populations for each trace of the 75.9kBq dataset. Green line is the population of $|0\rangle$, red line the population of $|1\rangle$, blue line the population of $|2\rangle$ and yellow line the population of $|3\rangle$.

An example of 2D histogram with corresponding fit for one of the data traces is shown in Fig. 6.7. As expected from the reset protocol, in most of the measurements the qubit is found in $|1\rangle$, there labeled as $|e\rangle$, which then has the largest cloud by far. In Fig. 6.8, instead, is shown the populations of the qubit per each trace.

6.3.2 Identifying the events

Measuring the qubit once in the ground state, or 0, does not automatically mean that a particle has hit the chip. Even if the waiting period, $5 \mu\text{s}$, is small compared to the decay time of the qubit, the probability to observe it decay within such time is non-zero. Across datasets, we observed that it ranged between 11.8 and 17.6% (see Table 6.3). In addition, other noise sources, like electric noise, vibrations [117, 174] and material relaxation at low temperatures [175, 176], can affect qubit behavior. As a consequence, we needed more stringent selection criteria.

As a first step, we required that the events we were selecting were compatible with a drop in the decay time, i.e. we were selecting events where we measured 0 multiple

times and not only once. As a trigger condition, then, we looked for sequences of four consecutive 0's in our data¹. This choice was motivated by the fact that the probability to measure this sequence in normal qubit conditions would be lower than 0.1%.

The trigger condition alone, however, is not sufficient to completely disentangle radiation-induced drops of the decay time from other noise mechanisms. It is true that observing four consecutive 0's in normal conditions is quite unlikely, but if we divide this probability by the sampling period (i.e. the time necessary to complete a full measurement cycle, including the cooldown period), obtaining a rough estimation of the expected rate of these events from spontaneous decay alone, we find a rate of a few events per seconds. When considering the expected contribution from ionizing radiation, estimated with simulations (Tables 6.1 and 6.2), we can see that even with the 161 kBq source, the most active one, we expect a rate of 0.43 events per second, much smaller than the rate coming from spontaneous decay. Particle impact events, then, would be completely obscured by noise events, especially when the qubit is operated deep-underground at LNGS, or exposed only to environmental radioactivity at FNAL.

The simplest solution to this problem could be increasing the consecutive number of 0's in the data stream necessary to trigger an event. Such a solution, though, is likely to discard also radiation-induced events, due to its stringent requirements: a sequence of ten 0's interrupted by a single 1 in the middle, for example, could be discarded despite it clearly being an event of interest. We opted then for a different approach.

We decided to keep the four consecutive 0's as the trigger condition, but then we added a second step of selection. We considered a window of 40 measurements (corresponding roughly to 2-3 ms given the sampling periods in our datasets) named *signal region*, starting five measurements before the trigger, and counted the number of 0's inside that window. We then retained only the events having more than a minimum number of 0's (N_{signal}^{min}).

Including in the signal window also the five points prior to the trigger ensured that the beginning of the event was also captured. The size of the signal window was optimized through an iterative process. Initially, a large window of 100 points was chosen. Then, since no events with more than 60 zeros were registered, the signal window was reduced to 60 points. With this new length, all the events contained 40 or less points. We concluded, then, that the longer events we were observing with the 100-point window were actually lasting shorter, but spurious 0's, arising from spontaneous decay once the qubit had returned to normal conditions, were affecting our count. Signal windows lasting less than 40 points were not used, since it was observed that they were actually

¹The entire analysis was repeated, for consistency, using three consecutive 0's as trigger condition, obtaining consistent results.

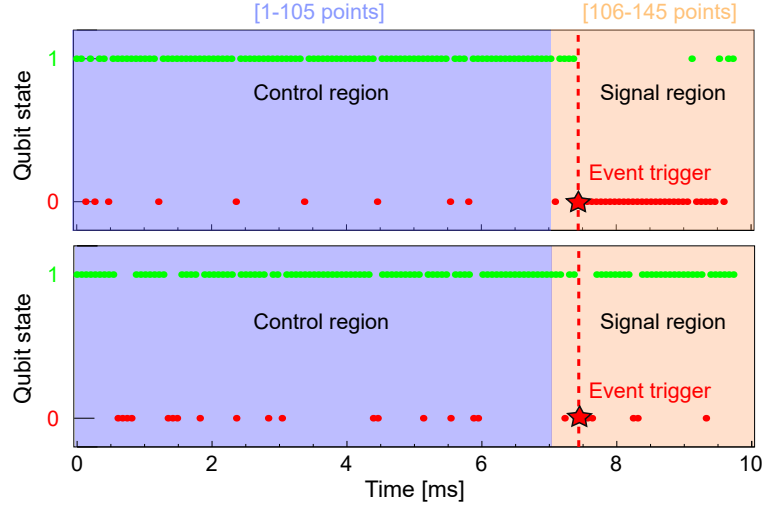


FIGURE 6.9: Two example events from the data, with the trigger position, the signal region and the control region highlighted. The event on top features a significant increase in $P(0)$ in the signal region compared to the control region, and so is selected. The event on bottom, instead, despite featuring four consecutive zeros, is characterized by a total number of zeros in the signal region comparable with what is observed in the control region, and is thus discarded.

shorter than some events. To avoid multiple triggers of the same event, a dead time of 35 points was applied after each trigger.

The choice of N_{signal}^{min} was determined by the expected rate of noise events from spontaneous decay. The probability to observe more than N_{signal}^{min} zeros in the signal window, coming only from the spontaneous decay of the qubit given its standard $P(0)$ value, can be computed by means of the binomial distribution. As for noise triggers, dividing this probability by the sampling period returns a rough estimation of the rate of noise events after this second step of selection. Since at LNGS, in the shielded configuration, we expect a rate of particle impacts in the chip of few 10^{-3} events/s, we chose for each dataset a value of N_{signal}^{min} that ensured a rate of events from spontaneous decay lower than 1×10^{-4} events/s. For a 40-point signal window, the values of N_{signal}^{min} resulted to be 21-22 for most of the datasets acquired. As an example, Fig.6.9 shows two events from our datasets, one with more than N_{signal}^{min} zeros in the signal window, thus passing the selection, and one with less, thus being discarded.

Finally, before computing the rate of events, we needed to make sure that this value was not poisoned by transient noisy periods of the qubit, similar to the ones observed in the experiment with the fluxonium. Because of that, we introduced a second window, that we named *control region*, including the 105 points that precede the signal window. Events with an anomalous number of 0's in the control window were flagged and discarded, as they could generate false triggers. Specifically, we computed the binomial probability of observing a given number of 0's in the control region and rejected events with a

probability lower than 1%. Since both an excess and a deficit of 0's can result in such low probabilities, we selected only events with a number of 0's in the control window between $N_{control}^{min}$ and $N_{control}^{max}$. The cut was intentionally chosen to be very loose, as its purpose was solely to remove events in which the qubit behavior was deviating significantly from the average. Both the control region and the signal region are highlighted in Fig. 6.9.

After the selection, the number of events registered in each dataset was divided by the corresponding live-time to estimate the rate of interactions. Uncertainties on the rates were estimated assuming a Poisson distribution of the events, so with a standard deviation equal to the square root of their number. Table 6.3 reports the computed rates of events in all the datasets, along with the acquisition time, the effective live-time after discarding traces with high state leakage, the sampling time, the mean $P(0)$ and the values of N_{signal}^{min} , $N_{control}^{min}$ and $N_{control}^{max}$. Datasets are reported following the same order in which they were acquired.

6.4 Results

6.4.1 Underground measurements

The first measurement was done underground at the INFN-LNGS, where we monitored the qubit both when fully shielded against radioactivity and when exposed to the Thorium sources. The results of the analysis are summarized in Fig. 6.10a and Fig. 6.10b.

Fig. 6.10a depicts the rate of observed events as a function of the number of zeros in the signal window. For all datasets, we observe a bell-shaped trend on the left side of the histogram. These events can be primarily attributed to spontaneous decays of the qubit: for these events, the probability to measure a certain number of zeros in the signal window, which is proportional to the rate, is determined by the binomial distribution, which follows that bell-like trend. After 20-21 zeros, however, the binomial bell rapidly decreases, and we can see that in the datasets with the sources there is a tail of long-lasting events which is absent in the dataset with the shielded qubit, pointing to a possible radiation origin.

Measuring with different radiation sources allowed us to confirm our hypothesis: if these long-lasting events are radiation-induced, their rate should increase with the activity of the source. This was indeed observed, as can be seen in Fig. 6.10b. The plot depicts the measured rate of events for each radiation exposure as a function of the expected rate of events, computed from simulations given the source activity. As expected, the rate of events increases linearly with the activity of the source, confirming that the tail

Dataset	Acquisition Time [min]	Live-time [min]	T_S [μ s]	$P(0)$ [%]	N_{signal}^{min}	$N_{control}^{min,max}$	Measured Rate [events/sec]
LNGS							
Shielded chip	736.7	210.0	73.6	11.8	19	6, 21	$(0.40 \pm 0.18) \times 10^{-3}$
44 kBq Th Source	620.7	114.0	67.6	14.5	21	8, 24	$(16.0 \pm 1.5) \times 10^{-3}$
76 kBq Th Source	124.1	46.3	67.6	13.8	21	7, 23	$(16.6 \pm 2.5) \times 10^{-3}$
125 kBq Th Source	62.1	23.7	67.6	15.5	22	9, 25	$(26.8 \pm 4.4) \times 10^{-3}$
161 kBq Th Source	62.1	62.1	67.6	14.9	21	8, 25	$(29.4 \pm 2.8) \times 10^{-3}$
FNAL							
60 μ s - run1	240.0	55.0	60.0	14.7	21	8, 23	$(7.9 \pm 1.5) \times 10^{-3}$
60 μ s - run2	240.0	180.0	60.0	15.8	22	9, 24	$(6.47 \pm 0.77) \times 10^{-3}$
60 μ s - run3	130.0	90.0	60.0	15.5	22	9, 24	$(7.9 \pm 1.0) \times 10^{-3}$
68 μ s - run1	270.4	270.4	67.6	14.2	21	8, 23	$(5.26 \pm 0.80) \times 10^{-3}$
68 μ s - run2	270.4	270.4	67.6	14.3	21	8, 24	$(5.73 \pm 0.59) \times 10^{-3}$
68 μ s - run3	270.4	270.4	67.6	13.5	20	7, 23	$(5.79 \pm 0.60) \times 10^{-3}$
74 μ s - run1	294.4	294.4	73.6	15.1	21	9, 25	$(5.32 \pm 0.55) \times 10^{-3}$
74 μ s - run2	294.4	294.4	73.6	14.5	21	8, 24	$(4.47 \pm 0.50) \times 10^{-3}$
74 μ s - run3	294.4	294.4	73.6	14.2	21	8, 23	$(4.98 \pm 0.53) \times 10^{-3}$
74 μ s - run4	294.4	294.4	73.6	15.0	21	8, 24	$(4.13 \pm 0.48) \times 10^{-3}$
50 μ s - run1	200.0	200.0	50.0	16.0	22	9, 26	$(8.33 \pm 0.83) \times 10^{-3}$
50 μ s - run2	200.0	200.0	50.0	16.3	22	9, 26	$(8.58 \pm 0.95) \times 10^{-3}$
50 μ s - run3	200.0	200.0	50.0	17.6	23	11, 28	$(9.08 \pm 0.87) \times 10^{-3}$
40 μ s - run1	159.2	159.2	39.8	15.8	22	9, 25	$(8.48 \pm 0.94) \times 10^{-3}$
40 μ s - run2	159.2	159.2	39.8	15.7	22	9, 25	$(8.58 \pm 0.95) \times 10^{-3}$
40 μ s - run3	159.2	159.2	39.8	15.0	22	8, 24	$(10.3 \pm 1.0) \times 10^{-3}$
55 μ s - run1	220.4	220.4	55.1	15.9	22	9, 26	$(7.26 \pm 0.74) \times 10^{-3}$
55 μ s - run2	220.4	220.4	55.1	16.2	22	9, 26	$(7.57 \pm 0.76) \times 10^{-3}$
55 μ s - run3	220.4	220.4	55.1	16.2	22	9, 26	$(7.26 \pm 0.74) \times 10^{-3}$

TABLE 6.3: List of the datasets collected at LNGS and at FNAL, along with their acquisition time, the effective live-time after removing traces with leakage, the sampling period, the measured probability to have $|0\rangle$ after the 5 μ s delay, the minimum number of zeros in the signal window to identify an event as radiation-induced, the minimum and maximum number of zeros in the control region required to keep the event, and the measured rate of events.

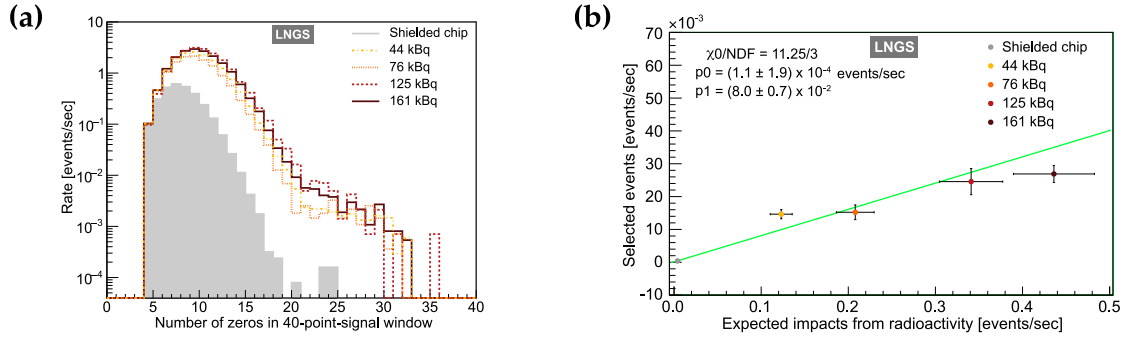


FIGURE 6.10: (a) Measured rate of events as a function of the number of zeros in the signal region, for the qubit in the shielded configuration and when exposed to the Thorium sources. (b) Rate of the measured events for the same datasets, vs the rate expected from simulations.

observed in the histograms is radiation-induced. This experiment also represents the first time that a superconducting qubit was operated as a γ detector, something that prior to that has only been hypothesized [177]. This result was not obvious, since γ -rays deposit less energy in the substrate compared to cosmic rays [131, 149].

Despite the positive outcome, one issue in our result requires more discussion. Fig. 6.10b, in fact, shows the rate of observed events as a function of the expected rate. The ratio of these events, which can also be extracted through the slope of the fit, can be thought as a *detection efficiency* of the qubit. This ratio, however, is just 8%. The main reason for that is the strong analysis cut imposed with $N_{\text{signal}}^{\text{min}}$: because of the predominance of spontaneous decay, only the longest-lasting events passed the selection, with the majority of the triggered ones being discarded. In this way, though, also shorter radiation-induced events may also have failed to pass the selection. This is also the reason why it was essential to compare datasets with similar $P(0)$ and populations. Fluctuations in these values would have resulted in different detection efficiencies and complicated the comparison between the rates.

Two complementary approaches can increase the detection efficiency. The first and most obvious one consists in choosing a qubit with higher T_1 : if $P(0)$ is lower, also the rate of events coming from spontaneous decay will drop, and it will be possible to choose a smaller $N_{\text{signal}}^{\text{min}}$. The second one, instead, consists of reducing the sampling period: in this way, a radiation-induced event lasting a few milliseconds will produce a higher number of zeros in the data stream, making it easier to identify. The latter approach has been successfully tested in the second part of the experiment, when the qubit was characterized above-ground at FNAL, and will be described in the next section.

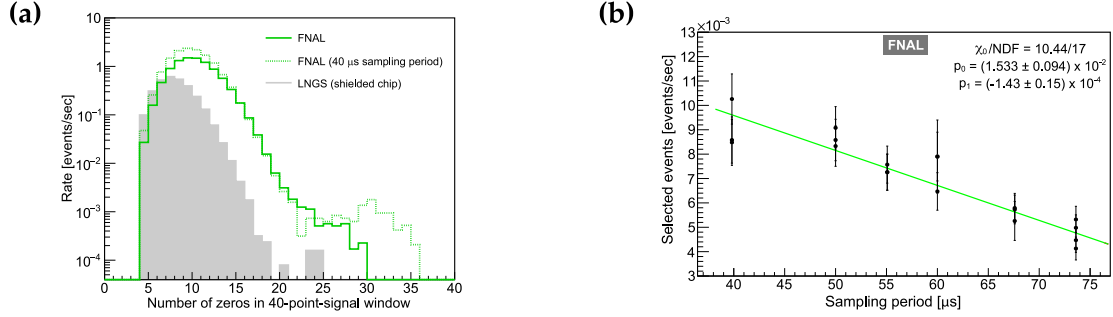


FIGURE 6.11: Results of the above-ground measurements at FNAL. (a) Histogram of the rate of detected events versus the number of zeros in the signal window, compared with what obtained at LNGS in the shielded configuration. The dotted line depicts what obtained in a different dataset with shorter sampling period: the radiation-induced tail is much more visible, due to the higher number of zeros produced by a millisecond-long quasiparticle poisoning produced by radiation impacts. (b) Rate of measured events as a function of the sampling period: the plot shows how the detection efficiency can be improved significantly by reducing the sampling period.

6.4.2 Above-ground measurements

Measurements at LNGS proved that it is possible to detect γ radiation with superconducting qubits. The following step was then to prove that, when operated above-ground with no radiation shielding, the measured rate increased by a factor 10, as expected from simulations. This was confirmed when we carried out the same measurement on the same qubit at FNAL.

The results of these measurements are summarized in Fig. 6.11a and Fig. 6.11b. The green line in Fig. 6.11a shows the rate of measured events as a function of the number of zeros in the signal window for the datasets with sampling period of 73.6 μ s, the same used in the shielded LNGS datasets. For comparison, the rate for the shielded LNGS dataset is also portrayed, in gray.

As expected, in the datasets taken at FNAL it is possible to observe a tail of long-lasting events that was not present when the qubit was shielded against ionizing radiation. These events are analogous to the ones observed when exposing the qubit to the Thorium sources, confirming their radiation origin.

This result conclusively demonstrates that with our protocol it is possible to detect particle impacts in the chip. The rate measured on the 73.6 μ s datasets at FNAL is $(4.68 \pm 0.26) \times 10^{-3}$ events/sec, approximately one order of magnitude higher than the one measured at LNGS in the shielded configuration, which was $(0.40 \pm 0.18) \times 10^{-3}$ events/sec. Comparing the rate above-ground with what expected from simulation, $(42 \pm 3) \times 10^{-3}$ events/sec, we can see that the detection efficiency is around 10%, similar to what obtained at LNGS.

Fig. 6.11a shows also (dotted line) the measured rate of events for the datasets with sampling period of $39.8 \mu\text{s}$. For these datasets, the tail of long-lasting events is much more evident compared to the $73.6 \mu\text{s}$ ones, and the measured rate doubled despite radiation exposure did not change (see Table 6.3). As anticipated in the previous Section, this effect can be ascribed entirely to the sampling period: radiation-induced events are lasting a few milliseconds regardless of the sampling period is, but if the sampling period is reduced, then events with the same duration will produce a longer stream of zeros in the data. At the same time, $P(0)$ remains similar across datasets, so the bulk of the histogram will not vary much, and $N_{\text{signal}}^{\text{min}}$ will stay approximately the same. As a consequence, radiation-induced events and events from spontaneous decay will overlap less in the histogram, and the detection efficiency will benefit from that.

To further confirm this result, we acquired several datasets with sampling times ranging from approximately $40 \mu\text{s}$ to approximately $73 \mu\text{s}$, as can be seen in Table 6.3. The analysis of these datasets showed that the measured rate of events decreases linearly with the sampling period, as can be seen in Fig. 6.11b. It is then possible to improve the detection efficiency of the qubit, even if T_1 remains unchanged, by reducing as much as possible the sampling period, provided that state leakage does not become significant. The sampling period, in fact, is determined mostly by the cooldown period after the state readout, as described in Section 6.2. That period, however, was added in the measurement protocol in order to prevent qubit heating and unwanted transitions to $|2\rangle$ and higher states, that would reduce its live-time. For this reason, datasets with sampling periods below $40 \mu\text{s}$ were not acquired. Still, reducing the sampling period to $40 \mu\text{s}$ was already enough to improve the detection efficiency by a factor 2 compared to the measurements at LNGS-INFN.

6.5 Summary and concluding remarks

The experiment on the transmon qubit described above allowed us to draw several important conclusions. The first is the confirmation that superconducting qubits are sensitive to ionizing radiation, as also shown in other studies. With the measurement and analysis scheme that we developed for this study, we were able to disentangle single particle impacts in the chip substrate, showing that the quasiparticles produced by these impacts recombine within a few milliseconds. This finding confirms our hypothesis regarding why we were unable to observe a statistically significant impact of ionizing radiation on the fluxonium chip: since the quasiparticle poisoning lasts only a few milliseconds, and since on a typical chip impacts occur only once every few seconds, a standard T_1

measurement, that requires several seconds to be completed, is not sensitive to such rare and fast fluctuations.

However, the long stream of zeros produced by particle impacts proves that the decay time of qubits drops well below the $5\ \mu\text{s}$ chosen as delay for our experiment, meaning that any computation performed when the chip is hit will end up failing. Such a low T_1 is also problematic for error correction algorithms, that require the error rate to be lower than a threshold to be effective. Considering also the results from Wilen et al. [131] and McEwen et al. [133] on radiation-induced correlated errors, we have to conclude that environmental radiation is an issue that must be addressed to get fully-functioning quantum computers, even if it is not the main decoherence source of superconducting qubits, otherwise any complex computation, lasting more than a few seconds, is likely to fail.

Another important achievement of this experiment is the demonstration that superconducting qubits are sensitive also to γ -rays and not only to cosmic rays. γ -rays deposit on average less energy than cosmic rays in the chip substrate, but their impacts are also much more frequent, making them a significant source that must be taken into account.

Finally, our experiment was the first to prove that superconducting qubits can be used as detectors for gamma radiation, something that prior to that was only hypothesized. The achieved performance, with an estimated detection efficiency of just 10%, is far from the over 99% efficiency of state-of-the-art particle detectors. Nonetheless, this result was achieved on a qubit that was not engineered for particle detection, with a detection protocol that may be further improved or even replaced by a more efficient one. This result should therefore be considered as a first step in that direction, with substantial R&D still necessary before achieving a fully developed particle detector. With their sensitivity to single quasiparticle tunnelings, and a typical binding energy for Cooper pairs of $O(\text{meV})$ and lower, superconducting qubits could serve as building blocks for detectors sensitive to energy deposits below 1 eV, making them of interest also for the community of particle physicists.

Overall, our experiment achieved several important results; however, many open questions remain, ranging from the estimation of the energy deposited by an impacting particle, and thus the minimum deposit required to induce decoherence, to the propagation of phonons through the substrate and the test of different mitigation strategies against correlated errors. The final part of this PhD was then devoted to improving and testing the room-temperature electronics used for our measurements, in order to enable the simultaneous drive and readout of multiple qubits, and is the subject of the next chapter.

Chapter 7

Detecting correlated errors with multiplexed control and readout

The results achieved with the study of the transmon qubit deepened our understanding of how ionizing radiation affects qubit coherence, but did not investigate the other major concern for quantum computing, which is correlated errors. At the same time, probing multiple qubits on the same chip would yield valuable insights into phonon propagation in the substrate, helping answer many questions arising from single-qubit studies. For example, understanding the impact of the ground plane could be useful for the design of future quantum-computation chips, given its role in trapping phonons produced in the substrate. From a particle detection perspective, instead, open questions span from estimating the energy deposited by the impact to improving the discrimination of noise events, which would be highly beneficial for detection efficiency, and so on.

However, studying correlated errors requires probing more than one qubit at the same time, meaning that the room-temperature electronics used in the experiment should be able to produce multiple pulses simultaneously, one for each qubit. This was not the case in the previous experiment, where the ZCU216 board used could handle only one qubit at a time.

The final part of my PhD, then, was devoted to overcoming this limitation. We developed a new firmware for the ZCU216 board in our lab, enabling the production of pulses capable of manipulating and reading the states of up to four qubits. We then updated all the software used for executing quantum algorithms and saving the output data, and tested everything on a new chip, developed by the SQMS Center. Finally, we improved the qubit reset based on the conditional π -pulse described in Chapter 6, developing a new version capable of resetting four qubits simultaneously. This new version is flexible enough to grant the possibility to set the qubits not only in $|0000\rangle$ or in $|1111\rangle$, that is, all

of them in the ground state or in the excited state, but also in other combinations, with some qubits in $|1\rangle$ and the others in $|0\rangle$. With this new reset protocol, an initial proof-of-principle measurement of correlated errors, carried out with an LED illuminating the qubits with periodical light pulses, was successfully completed, paving the way for future research in that field.

7.1 Generation and readout of multiplexed pulses

7.1.1 The QICK firmware for the control of RFSoc boards

The generation of microwave pulses by the ZCU216 RFSoc board, their timing, and the demodulation and readout of the returning signals (such as readout pulses) are controlled through the Quantum Instrumentation Control Kit (QICK) [172]. QICK consists of a firmware for the RFSoc board and a Python interface for controlling it from a PC.

The QICK firmware comprises several elements, or blocks, that have to deal with clocking, data flow, user interface and other functionalities, but three of its blocks are of particular interest in quantum computing experiments: signal generators, readouts and the *timed-processor*, or *tProcessor*.

The latter allows the implementation of quantum algorithms, controlling the generation and readout of pulses along with their timing. It features eight output channels and one input channel¹, and can be programmed using a custom assembly language. The instructions can be executed at specific times thanks to an internal clock and include: looping, conditional branching, memory access and triggering the generation and readout of signals.

The signal generator block, instead, deals with actually producing the RF pulses at the time specified by the tProcessor. QICK firmware versions usually feature more than one generator, depending on the requirements of the experiment (at least two in typical experiments, one for drive and the other for readout pulses). The generators implemented in the standard version of the QICK firmware are capable of producing pulses with only one carrier frequency, or tone, meaning that only one qubit can be operated at a time.

Signals are produced with a *Direct Digital Synthesizer*, or DDS, and then multiplied by the pulse envelope before being digitally upconverted to the required frequency. Pulse envelopes are saved in the FPGA memory in the form of I and Q arrays, to allow fast

¹A new, more flexible version of the tProcessor has been released only recently, so it was not used for the work of this thesis.

pulse generation. Upconversion is achieved by mixing the resulting pulse with a tone produced by the DDS itself. After all these steps, the pulse is amplified to the desired amplitude and finally sent to the Digital-to-Analog Converter, or DAC.

An important feature of the signal generator is *phase coherence*, meaning that pulses with the same frequency, even if separated in time, have a fixed relative phase between them, that can be set by the user. This feature has several implications for qubit experiments, from the orientation of the $I - Q$ clouds in a fidelity measurement, which is determined by the phase of the readout pulses and is fundamental for the correct functioning of the qubit reset described in Chapter 6, to the rotation axis on the Bloch sphere when driving the qubit, which depends on the phase of the drive pulse (see Section 2.1.2).

The readout block follows a process which, at least in its first part, reverses what is done by the signal generator. The incoming signal is first digitized by an Analog-to-Digital Converter, or ADC, and then digitally downconverted by mixing it with the tone produced by a DDS. After the mixing, the signal undergoes a low-pass filter to select only its low-frequency component and is then decimated by 8. In this way only the pulse envelope is saved, removing the signal oscillations associated with its carrier frequency. After that, the values of the signal components over time are averaged, returning a pair of single $I - Q$ values. These values can be saved in the board memory, in order to be accessible by the user, and also sent to the tProcessor for their use in conditional statements (like the instruction for the conditional π -pulse in the reset protocol). The entire readout process is analogous to the one described in Section 2.2.2.

7.1.2 Multiplexed signal generators and readouts

Firmware blocks for the generation and readout of multiplexed signals have been developed and are available in the QICK source code [178], but need to be implemented in the firmware in order to control and measure multiple qubits simultaneously. In our work, we modified the QICK firmware by including these blocks, creating a custom version that features two multiplexed signal generators, one for drive pulses and the other for readout pulses, and one multiplexed readout.

The multiplexed signal generator is conceptually identical to the single-tone one: each tone is produced by an independent DDS, multiplied by the pulse envelope and upconverted to the desired frequency. The resulting single-tone signals, however, are summed together before being sent to the DAC. The multiplexed pulse produced can be seen as the sum of the outputs of several standard signal generators.

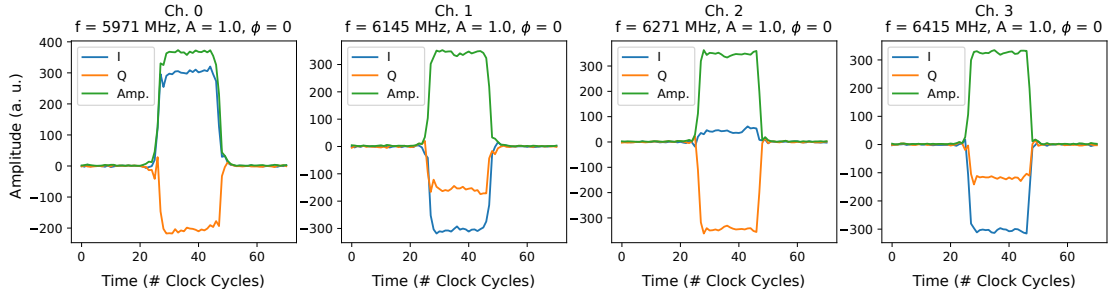


FIGURE 7.1: Demodulated I and Q values, and the signal amplitude, in arbitrary units, reported as a function of time for a four-tone pulse produced by the ZCU216 with the custom QUICK firmware. Time is expressed in terms of clock cycles. On top of each plot, A denotes the relative amplitude of the corresponding tone (ranging from 0 to 1) and ϕ is the phase.

An important caveat that must be taken into account when designing an experiment is that the maximum amplitude of each tone is only a fraction of that of a single-tone pulse. This is because the produced pulse undergoes the same DAC used for single-tone pulses. As a consequence, in the case of multiplexed pulses the maximum amplitude of the signal has to be split between its tones.

Concerning the multiplexed readout, this can also be seen as several single-tone readout blocks. The signal undergoes a *Polyphase Filter Bank* (PFB) [179], whose effect is to split it into separate channels with different frequency bands (analogous to applying in each channel a different band-pass filter). Each tone then undergoes the same steps of single-tone readouts, being first downconverted and then decimated and averaged to extract I and Q .

Fig. 7.1 shows an example of a multiplexed RF pulse with four different tones. The multiplexed generators employed in our custom firmware were able to generate signals with up to eight different tones. The readout, however, was able to demodulate signals with a maximum of four tones, limiting the number of qubits we could simultaneously operate. This limitation was due only to a lack of memory available on the board. The firmware developed, however, can be optimized further in order to use the available memory more efficiently, and new versions, capable of controlling and reading out eight qubits simultaneously, are likely to be developed in the future.

7.1.3 Test on real qubits

The new firmware was tested on a chip with eight transmon qubits developed by the SQMS Center. The geometry and fabrication of this chip is similar to the one described in Chapter 6, with the main difference being that the Niobium qubits deposited in this

one have no capping. The test consisted in a full characterization of four qubits and in the optimization of the corresponding readout fidelities.

Qubit characterization started with a spectroscopy experiment, probing the chip with signals at different frequencies, to determine the resonant frequencies of readout resonators and the transition frequencies of the qubits. Resonator spectroscopy, that can be seen in Fig. 7.2a, consisted of sweeping the frequencies of the readout tones to reconstruct resonances. Because of the low temperature of the dilution refrigerator, spontaneous excitations of the qubits are rare, so the positions of the resonances correspond to the qubits being in $|0\rangle$.

Qubit spectroscopy, for determining the transition frequencies of the qubits, is slightly different. This measurement consists of sending a strong drive pulse, to excite the qubit, followed by a readout pulse to check if the qubit was actually excited. Since the frequencies of the readout pulse were tuned to be at the resonance positions with the qubit in $|0\rangle$, if the drive pulse is effective, the resonance will shift, and the returning readout amplitudes will be higher. Following this idea, the frequencies of the drive tones are swept until a peak in the readout response is found, meaning that the resonance has shifted and, hence, the qubit has been excited. Fig. 7.2b shows the outcome of the qubit spectroscopy experiment, where the amplitudes of the readout tones are plotted as a function of the frequencies of the drive tones. The peaks on the plots determine ω_{01} of the qubits.

The next step was the determination of Rabi frequencies, which are fundamental for coherently manipulating the qubits' states. In Section 2.1.1 it was shown that Rabi oscillations can be observed by varying both the amplitude and the length of drive pulses. In the case of multiplexed signals, however, all the tones are generated for the same duration. Since different qubits have different Rabi frequencies, we decided to keep the pulse length fixed and sweep the amplitudes of the drive tones to observe oscillations. The result is portrayed in Fig. 7.2c.

Having determined Rabi frequencies, the drive pulse is working, but additional fine-tuning is still required. Qubit spectroscopy, in fact, is a necessary step to perform a Rabi experiment, but its estimation of the drive frequency is not precise. Sending strong drive pulses like the ones used in that experiment, in fact, widens the qubit resonance, since also off-resonant pulses can excite the qubits due to the sidebands in their Fourier transform. A more refined estimation of ω_{01} was then obtained by performing a Ramsey fringe experiment (see Fig. 7.3a), as described in Section 2.1.3. The experiment was performed by sweeping both the drive frequency and the time delay between the two $\pi/2$ -pulses. For each channel, the transition frequency was determined by looking at the central fringe in the interference plot.

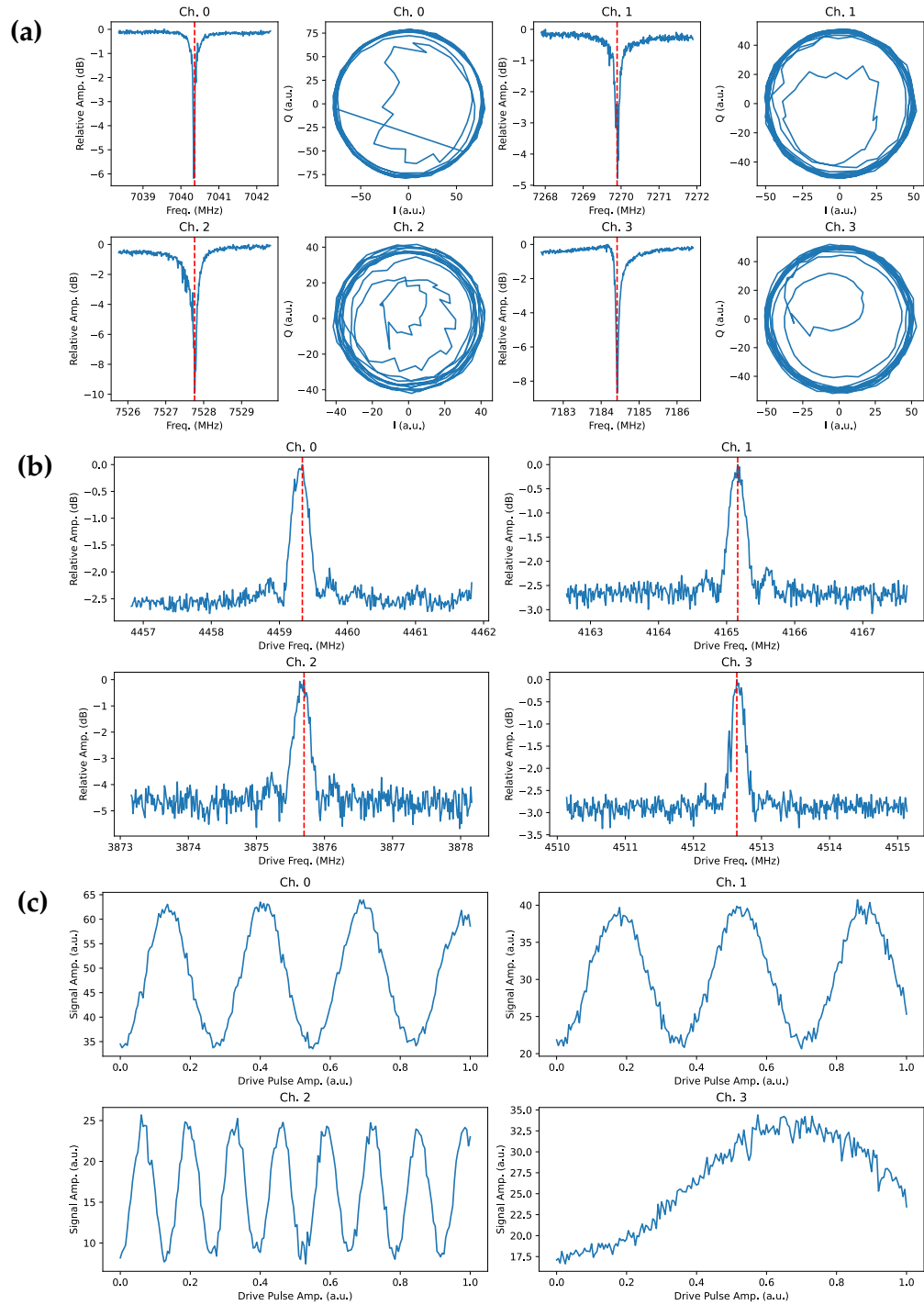


FIGURE 7.2: First part of the simultaneous characterization of four qubits with the custom QICK firmware. (a) Spectroscopy of the readout resonators to determine the central frequency. For each channel, the transmitted signal is shown both in the frequency-amplitude plot (left) and in the $I - Q$ plot (right) (b) Qubit spectroscopy to determine the transition frequency. Here, the amplitude of the measured readout pulse is plotted as a function of the drive pulse frequency. When the drive pulse is resonant with the qubit and excites it, the resonance of the readout resonator shifts and the amplitude of the readout pulse changes. (c) Rabi oscillations as a function of the tone amplitudes for 0.5- μ s-long drive pulses. The different Rabi frequencies are a consequence of the different coupling of the qubits with the feedline.

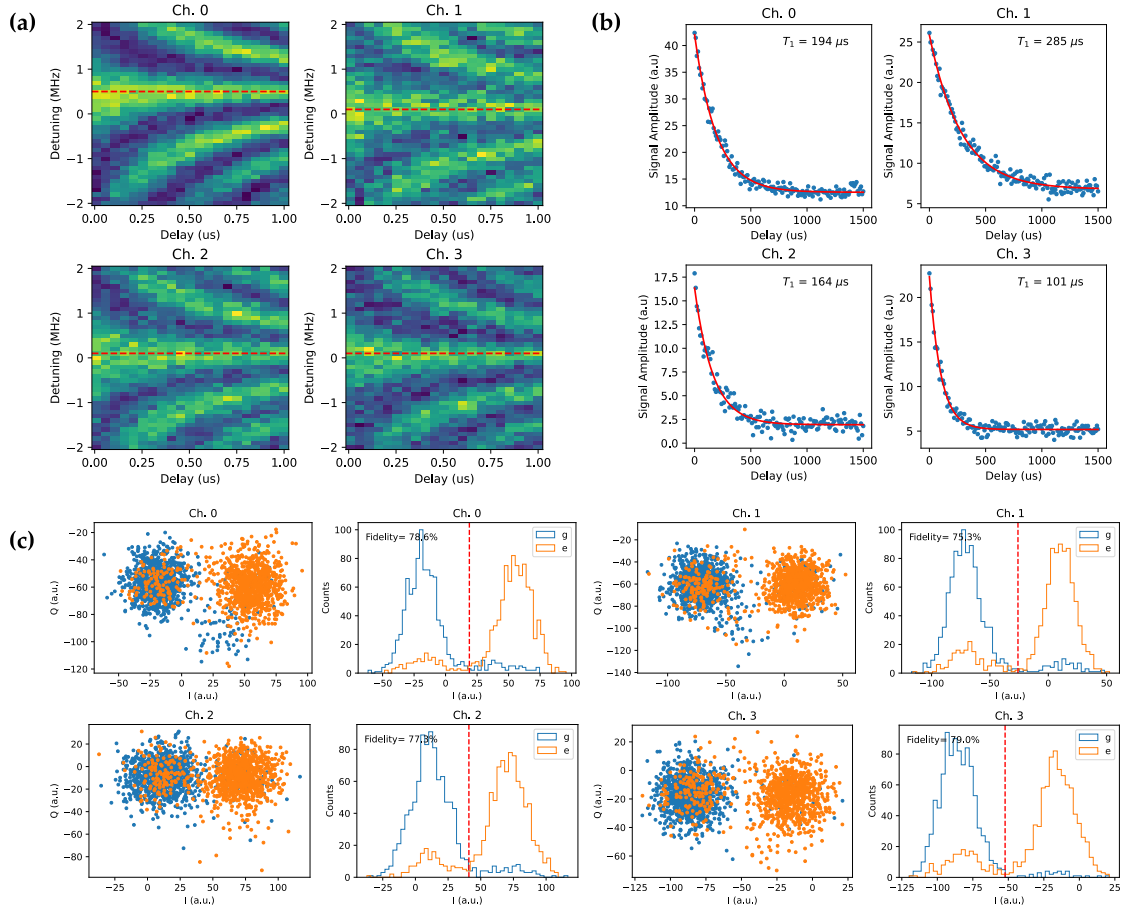


FIGURE 7.3: Second part of the simultaneous characterization of four qubits with the custom QICK firmware. (a) Ramsey fringe experiment (the detuning is computed from the transition frequency obtained with qubit spectroscopy). (b) T_1 measurement. (c) Fidelity measurements, with $I - Q$ clouds and their projections onto the I axis. The dotted red line indicates the optimized threshold for state discrimination. As can be seen from the plots, clouds are well separated, with the fidelity being limited mostly by residual $|g\rangle$ population after the π -pulse.

With an optimized drive pulse, it was finally possible to determine the two quantities of interest for our measurements: T_1 and the readout fidelity. T_1 was measured as in Section 2.3.1, by preparing the qubit in $|1\rangle$ and measuring the probability to have it still in $|1\rangle$ after a time delay. For all qubits we found T_1 values of hundreds of μ s, as can be seen in Fig. 7.3b, compatible with the values measured when they were characterized without multiplexing.

The readout fidelity, instead, was optimized to maximize the discrimination power along the I axis in the $I - Q$ plot (see Fig. 7.3c). This was done by sweeping over all the parameters of the readout pulse, so amplitude, frequency and duration, and looking for the combination that yielded the best result. As for the drive pulse in the Rabi experiment, also for the readout pulse all the tones must have the same duration. However, the pulse length yielding the maximum fidelity for one qubit may not be the most

suitable one for another. For our test, we decided to choose a pulse length that was yielding a good fidelity value on all channels, even though higher fidelity could be achieved for some qubits. Nonetheless, we still managed to achieve a measurement fidelity, i.e. $\mathcal{F} = 1 - P(1|0) - P(0|1)$, between 75 and 80%, consistent with measurements obtained from single-qubit operations.

Overall, this test demonstrated that the multiplexed firmware operates correctly without introducing additional noise into the measurements, with qubit characterization results consistent with those obtained using single-tone pulses.

7.2 Development of the multiplexed qubit reset

7.2.1 The conditional drive pulse

Having developed a multiplexed firmware for our board, the next step was the development of a new reset protocol, capable of resetting multiple qubits simultaneously. As for the single-qubit one, described in Section 6.2, we wanted a protocol that could be applied to any type of superconducting qubit, without relying on specific hardware features. Following this approach, the new reset protocol that we developed shared with the single-qubit one the same working principle: a readout pulse has to be sent first, to see which qubits are in $|0\rangle$ and which ones are in $|1\rangle$; after that, a conditional π -pulse excites the qubits that are found in $|0\rangle$. However, the multiplexed reset faced a complication that was not present in the single-qubit one: the fact that different qubits could be in different states. In this scenario, a standard multiplexed π -pulse would have excited the qubits found in $|0\rangle$, but also sent to $|0\rangle$ the qubits that were already in $|1\rangle$.

The solution consisted in implementing, in our algorithm, four different conditions, one for each qubit that had to be reset (see Fig. 7.4a). The reset starts with the amplitudes of all tones of the π -pulse set to 0. Then, for each qubit the algorithm reads the corresponding readout tone, adjusting the amplitude of the drive tone only if the reset condition gets satisfied.

As for the single-qubit reset, these conditions could be implemented only if one numeric value was enough to identify the qubit state. The phases of the readout tones were then adjusted in such a way that the $I - Q$ clouds of each qubit, corresponding to the $|0\rangle$ and the $|1\rangle$ state, were centered on the same Q value. In this way, as in Chapter 5 and Chapter 6, separation between the clouds would be maximized along the I axis, allowing the determination of the qubit state based only on the demodulated I value. The reset

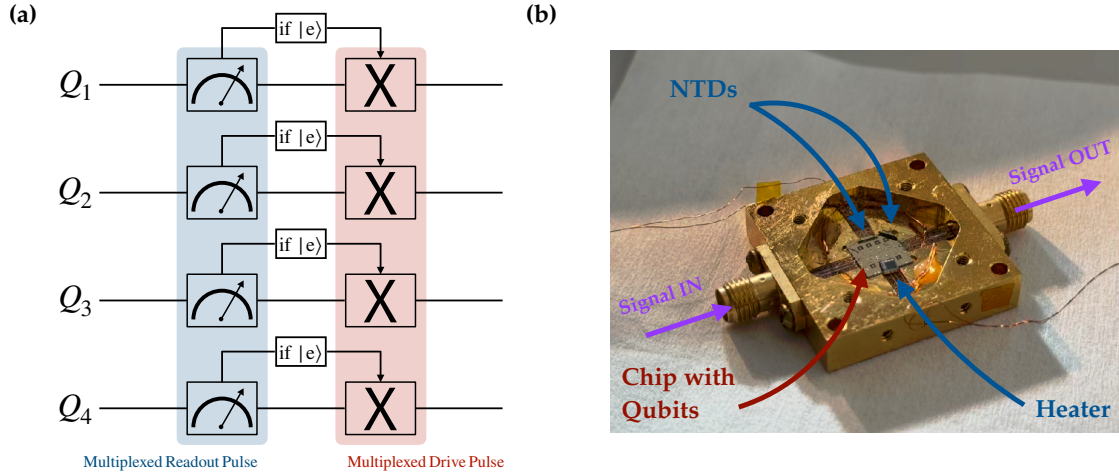


FIGURE 7.4: (a) Scheme of the multiplexed reset protocol for the simultaneous reset of four qubits to $|1111\rangle$. By changing the conditions for the π -pulse tones, it is possible also to reset the qubits in different states. (b) The chip used for the test of the multiplexed qubit reset and the first measurements of correlated errors. Two NTDs are also glued on the chip, for the coincident detection of particle interactions in the chip, and a heater for the injection of thermal phonons.

condition is then implemented by comparing the readout I value with a threshold set by the user, corresponding to a value in between the two clouds.

The sign of the comparison determines whether the qubit has to be reset in $|0\rangle$ or in $|1\rangle$. Having four, independent conditions, then, grants much higher flexibility than typical multiplexed reset protocols, that reset all qubits in $|0\rangle$ and, only afterwards, excited the desired qubits to $|1\rangle$ [134, 180–183]. In our protocol, instead, each qubit can be equally set either in $|0\rangle$ or in $|1\rangle$ simultaneously with the others, without being affected by the other reset conditions. As a consequence, qubits can be prepared not only in $|0000\rangle$ (typical final state for most of the reset protocols) or $|1111\rangle$ (the desired state for our experiment on radiation-induced errors), but also in $|0001\rangle$, $|0100\rangle$, $|1100\rangle$ and so on, granting a flexibility that is not common in these algorithms.

7.2.2 Test and fidelity measurements

The chip used for the test of the new firmware would have been perfect for a correlated error measurement based on the protocol described in Section 6.2: the high T_1 values of its qubits would have granted much lower noise events compared to the transmon used in Chapter 6. However, the readout resonators mounted were unsuitable for the multiplexed qubit reset just developed. This was due to their very long *ringdown time*: long ringdown times mean that the conditional π -pulse being sent right after the readout would interfere with the previous signal inside the resonator, compromising the qubit drive. The only way to avoid it would be to introduce a delay of a few microseconds

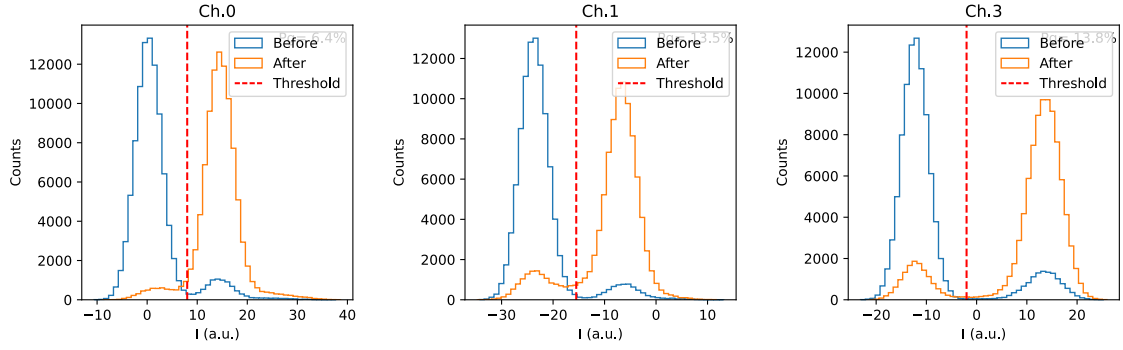


FIGURE 7.5: I values of qubits' readout before (blue) and after (orange) the reset, for the reset in $|111\rangle$. The dotted red line is the threshold for state discrimination. Here Channel 2 is not portrayed because for that qubit it was not possible to determine its transition frequency, so the reset algorithm was not effective.

between the readout and the π -pulse, but that would significantly slow down the qubit reset and make it more unreliable, since the longer time passes between the two pulses the higher is the probability that the qubit state changes meanwhile.

For this reason, we decided to employ a different chip for the test of the multiplexed qubit reset. The chip chosen was also developed by the SQMS Center, following a design and a geometry analogous to the one used for the single-qubit gamma detection experiment in Chapter 6. The materials chosen were also similar, with eight niobium qubits deposited on a sapphire substrate, incorporating Al/AlO_x/Al Josephson junctions. The main difference with the other chip is that the 10-nm-thick capping of niobium pads is made of tantalum instead of gold [89]. The readout resonators of this chip have ringdown times of $O(100\text{ ns})$, much faster than the ones of the other chip, making them suitable for the implementation of the multiplexed reset.

Once the test of the reset protocol was completed, the chip was intended to be used for the measurement of correlated errors. To better identify particle-induced events, two germanium *Neutron Transmutation Doped* (NTD) thermistors [184] were glued on the chip. They are made of doped germanium wafers with a resistivity strongly depending on their temperature, and are used in cryogenic particle physics experiments for the detection of particle interactions through the phonons produced. On the chip we also glued a heater, i.e. a resistor, for the injection of thermal phonons through small electric signals.

The glue used for this process was the two-component epoxy Araldite[®] Rapid, particularly convenient for cryogenic experiments because of its low-temperature performance and high radiopurity. Unfortunately, being an amorphous material, it has a high concentration of two-level systems. This ended up degrading severely the qubits on the chip: both T_1 and T_2 dropped significantly, with T_1 being of a few tens of μs and T_2 being $O(\mu\text{s})$, the readout fidelity also worsened and only for three qubits out of eight a full

		Measured State							
		000	001	010	011	100	101	110	111
Prepared State	000	0.76911	0.11038	0.04382	0.00689	0.05734	0.00839	0.0035	0.00057
	001	0.22644	0.59635	0.0202	0.08408	0.01688	0.0471	0.00178	0.00717
	010	0.04178	0.01599	0.47403	0.19151	0.01403	0.00704	0.16804	0.08758
	011	0.01207	0.07064	0.11572	0.75248	0.00066	0.00382	0.00555	0.03906
	100	0.06156	0.00548	0.00654	0.00051	0.76135	0.06626	0.09133	0.00697
	101	0.01819	0.05088	0.00014	0.00094	0.24163	0.67418	0.00231	0.01173
	110	0.00499	0.00072	0.04575	0.00694	0.07502	0.01116	0.74733	0.10809
	111	0.00139	0.00737	0.00739	0.04771	0.01904	0.10728	0.11013	0.69969

FIGURE 7.6: Error matrix for the multiplexed qubit reset. Each row corresponds to a different three-qubit state prepared with the reset protocol. The columns, instead, correspond to the three-qubit state that was actually measured after the reset. In each cell is reported the probability to measure a three-qubit state after the reset given the one that the protocol was required to prepare. Color of the cells gets darker as the probability increases.

characterization was achieved, while for others a good manipulation and readout of their state was impossible. Due to this limitation, we had no other choice than restricting the test of the multiplexed qubit reset to the three functioning qubits left, which in the processing of the signals produced by the ZCU216 have been labelled Ch. 0, Ch. 1 and Ch. 3. Ch. 2 was associated to a fourth qubit, for which we could not observe Rabi oscillations, making it impossible to coherently manipulate its state.

The reset protocol was tested by preparing the qubits in all possible three-qubit states. Fig. 7.5 shows, as an example, the outcome of the reset to the $|111\rangle$ state. In the plots, the blue histogram is the outcome of the readout pulses prior to the conditional π -pulse,

while the orange one is the outcome of the measurements done right after the reset. As it can be seen, in most of the cases qubits get prepared correctly, but in a small fraction of measurements some of them are found in $|0\rangle$ instead of $|1\rangle$. These fractions correspond to 6.4% for Ch. 0, 13.5% for Ch. 1 and 13.8% for Ch. 3, and can be ascribed to two main reasons: errors in the identification of the initial states, which are normal when the readout fidelity is not 100%, and state transitions during the readout after the reset, which are a consequence of the low T_1 of the qubits. Because of that, the gluing of the NTD and of the heater on that chip ended up degrading the performance of the qubit reset, even if still allowed for correlated error studies (see next Section).

Since the preparation of each qubit came with an error, the probability to initialize all of them in the correct state ended up being the product of the probabilities for the correct initialization of each qubit. For the preparation in $|111\rangle$, that probability corresponded to roughly 70%. An error matrix with the probability to get each outcome given the state required is depicted in Fig. 7.6. From the table, it is possible to see that in most of the cases the qubits were initialized to the correct states, with the highest probabilities along the diagonal. However, the errors in single-qubit preparation discussed earlier limited the final outcome, with the probability for the correct state to be measured after the preparation never exceeded 77%.

7.3 First measurement of correlated errors

With the reset protocol working, a first measurement of correlated errors was performed on the same chip described in the last section. Errors were induced by illuminating the chip with an LED, with 15 ms light pulses applied at one-second intervals. The total energy deposited in the chip by a light pulse was estimated with the NTD detector to be around 1.5 MeV.

7.3.1 Event triggering and selection

The event triggering was analogous to the method previously employed for the detection of gamma radiation, described in Chapter 6. As for that experiment, we used as trigger condition four consecutive zeros in the data stream. To flag an event, the trigger condition had to be satisfied by at least one qubit. For each trigger, both a *signal* region and a *control* region were acquired, defined as the 300 points following and the 100 points preceding the trigger, respectively. For this analysis the signal window was chosen to be much longer than the one used for the experiment with the gamma sources (in which it was only 40 points) because of the long duration of light pulses. As an example,

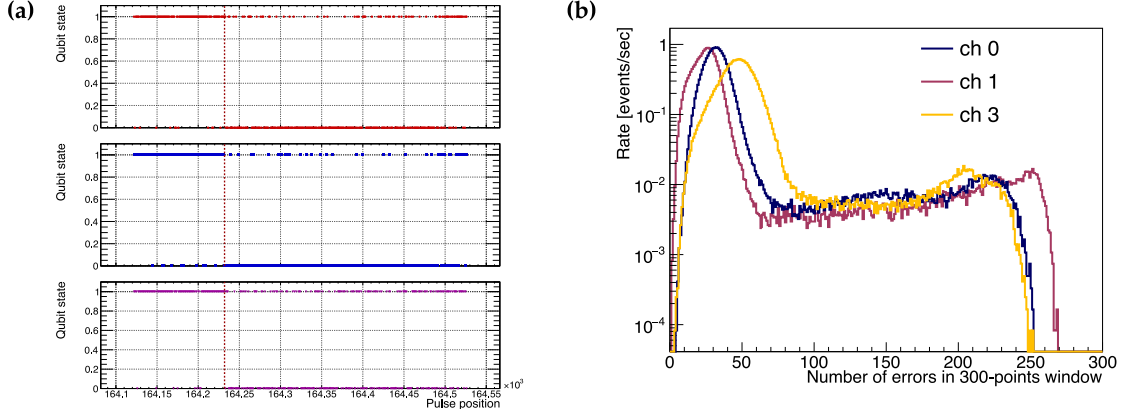


FIGURE 7.7: (a) Example of a correlated-decay event in three qubits generated from an LED light pulse. The x axis values represent the position of the measurements in the dataset partial. (b) Histograms of the number of 0's of the triggered events in the three qubits, where the tails of LED-induced events lasting hundreds of measurements are clearly visible.

Fig. 7.7a shows a simultaneous event on the three qubits caused by a light pulse. Fig. 7.7b, instead, shows the histograms of the number of zeros in the signal regions for the three qubits: after the initial bulk of events, that can be ascribed to spontaneous decay and follows a binomial distribution, the histograms feature a tail of long-lasting events reaching almost 300 zeros, which, in contrast, arises from the LED.

For the event selection we modified the strategy adopted in Chapter 6. In the analysis of the single-transmon experiment, in fact, we required a fixed amount of zeros in the signal region as the minimum to select an event, that we named N_{signal}^{min} . The threshold was selected based on the qubit's average $P(0)$, ensuring that the probability of observing this many zeros in the signal region, assuming only $P(0)$, was below a predetermined value. By scaling that probability by the sampling period of the qubit, the expected rate of these background events should be one order of magnitude lower than the expected rate of radiation-induced events at LNGS in the fully-shielded configuration. Such approach had the advantage of keeping the noise rejection efficiency constant across the data, but did not account for fast fluctuations of $P(0)$. As can be seen in Table 6.3, in fact, out of twenty-four datasets acquired at LNGS and FNAL, only three of them have a live-time shorter than one hour. But T_1 scans of the same transmon (see Fig. 6.1c) showed that the decay time of the qubit, and hence $P(0)$, fluctuates on timescales much shorter than that, minutes or even seconds.

The selection strategy in the experiment with the LED overcame this limitation by substituting the fixed N_{signal}^{min} with a cut chosen event by event. The starting point of this method is the number of zeros in the control region, defined as $k_{control}$. This number follows a binomial distribution, as in the analysis described in Chapter 6, but this time, since we account for fluctuations of qubit's T_1 , the probability $P(0)$ is unknown. We

compute, then, a posterior distribution for $P(0)$ based on the size of the control region and the number of zeros measured within it. To improve readability, from now on $P(0)$ will be denoted as p , and its posterior distribution as $P(p|k_{\text{control}}, N_{\text{control}})$.

Having computed the posterior distribution for p , the next step is to determine the probability that, given a minimum k_{signal} for the selection of an event, a trigger uncorrelated with the LED light passes the selection. In other words, the probability $P_B(k_{\text{signal}})$ to observe *at least* k_{signal} zeros in the signal region due only to the unperturbed behavior of the qubit. Given a value of p , this probability is given by the binomial distribution $B(k|N_{\text{signal}}, p)$:

$$P_B(k_{\text{signal}}|N_{\text{signal}}, k_{\text{control}}, N_{\text{control}}, p) = \sum_{k=k_{\text{signal}}}^{N_{\text{signal}}} B(k|N_{\text{signal}}, p) \quad (7.1)$$

Since p is unknown, however, we need to integrate over all its values, weighting each of them by the posterior probability found earlier:

$$\begin{aligned} P_B(k_{\text{signal}}|N_{\text{signal}}, k_{\text{control}}, N_{\text{control}}) \\ = \int_0^1 \sum_{k=k_{\text{signal}}}^{N_{\text{signal}}} B(k|N_{\text{signal}}, p) \cdot P(p|k_{\text{control}}, N_{\text{control}}) dp \end{aligned} \quad (7.2)$$

The result is that, for each event, we have an estimation of the probability that a noise event passes the selection given a cut. For the choice of the cut, then, we followed an approach analogous to the one described in Chapter 6. Since the LED was emitting a pulse every second, we wanted to be sure that the rate of noise events remained much lower than one event per second. For each event the minimum k_{signal} required is set in such a way that, given the posterior distribution of p , the rate of noise events does not exceed 10^{-1} events/sec, one order of magnitude lower than the rate of LED events. As for the single-qubit experiment, the rate of noise events was roughly estimated by dividing the probability that a noise event would pass the selection, given the chosen cut, by the sampling rate of the qubits, which was $60.7 \mu\text{s}$.

7.3.2 Measured rates

For this experiment, three datasets of different duration were acquired. The data streams from the three qubits were analyzed both individually and simultaneously, looking for coincident events. Results are summarized, respectively, in Table 7.1 and Table 7.2.

Dataset	Acquisition Time [min]	Rate [events/sec]		
		Ch. 0	Ch. 1	Ch. 3
Dataset 1	20	0.99 ± 0.03	1.01 ± 0.03	0.95 ± 0.03
Dataset 2	202	0.978 ± 0.009	1.005 ± 0.009	0.892 ± 0.009
Dataset 3	506	0.976 ± 0.006	0.997 ± 0.006	0.904 ± 0.005
Average		0.977 ± 0.005	1.000 ± 0.005	0.902 ± 0.004

TABLE 7.1: Acquisition times and rate of events in the three channels for the datasets acquired.

Dataset	Acquisition Time [min]	Rate [events/sec]
Dataset 1	20	0.94 ± 0.03
Dataset 2	202	0.908 ± 0.009
Dataset 3	506	0.899 ± 0.005
Average		0.902 ± 0.004

TABLE 7.2: Acquisition times and rate of coincident events in the three channels for the datasets acquired.

The first noticeable result is the high detection efficiency of LED pulses by the three qubits. For all the datasets, the measured rate of events in Ch. 1 is compatible with 1 event/sec, indicating a 100% efficiency. The performance of Ch. 0 is just slightly worse, while Ch. 3, the worst performing one due to higher values of p , still reached 90% efficiency.

Results are very similar when considering coincident events. In that case, every time an event was triggered in a channel, the control and signal regions were acquired for all. If all the channels, then, passed the selection cut, the event was registered. Also in this case, measured rates were close to 1 event/sec in all the datasets, with an average rate of correlated errors of 0.902 ± 0.004 events/sec, pointing to 90% detection efficiency. This value is limited mostly by Ch. 3, due to its lower efficiency.

These rates are a significant improvement compared to what achieved with the gamma sources, when the detection efficiency reached only 8%. However, there are two important differences between the two experiments that have to be taken into account for a proper interpretation of these results. The first one is the time duration of the events. In the case of gamma interactions in the substrate, the quasiparticles produced recombined within two or three milliseconds, meaning that the stream of zeros in the data lasted no longer than 30 or 35 measurements. In this experiment, instead, each LED pulse was illuminating the qubits for 15 ms, producing events with 250 up to 300 zeros. This high amount of zeros significantly simplifies the event selection, since in normal conditions

the qubit would not produce a noise event lasting so long. The high rate of LED events led also to the choice of a selection cut much less stringent than the one used in the experiment with the gamma sources, where for the shielded dataset the rate expected from simulations was $O(10^{-3})$ events/sec.

The second consideration is that, since the LED was placed right above the qubits, the emitted light was illuminating the superconductor. This means that the energy deposited was breaking Cooper pairs without undergoing phonon mediation, and it can be expected, in such scenario, that the detection efficiency is higher.

Overall, however, this result proved that simultaneous qubit decays can indeed be detected with multiplexed firmware and reset, and is an important step towards the study of radiation-induced correlated errors.

7.4 Summary and discussion of the results

The tests of multiplexed firmware, along with the successful operation of the multiplexed qubit reset for the detection of correlated events, allow us to draw some important conclusion for future research in this field.

The first, and most obvious, is that the electronics and software employed in our experiment are capable of simultaneously probing multiple qubits without degrading the signal-to-noise ratio of the pulses produced. This was a fundamental step towards the study of correlated errors, since the multiplexed reset protocol would not function in a different scenario: a higher electronic noise would result in greater overlap of the clouds and lower fidelity, while a strong cross-talk between channels would undermine the correct reset of the qubits.

The test of the multiplexed reset protocol, instead, showed the feasibility of a reset system capable of operating also on fixed-frequency qubits, and that can be applied to a wide range of superconducting qubit designs.

Finally, the successful detection of correlated qubit errors induced by the LED paved the way for future studies on the impact of radiation on multi-qubit chips, a field with profound implications both in quantum computing and in particle physics. As better understanding of this phenomenon will lead to radiation-resistant designs and algorithms, in fact, the sensitivity of qubits to particle interactions can also be the starting point for the development of a new type of particle detector. The operation of a transmon qubit as a gamma detector, that was described in Chapter 6, along with the detection of LED light described in the previous section, proved that superconducting qubits can be used

for this scope, but many questions remain open, and research in this field is still at an early stage.

Conclusion

As concerns grow about the negative impact that environmental radioactivity can have on quantum computers, especially next-generation devices, this thesis sheds light on how it affects the decoherence of superconducting qubits. While most of the previous studies [131, 133] focused on cosmic rays, especially muons, our Monte Carlo simulations showed that the majority of particle interactions in the chip substrate are due to environmental γ -rays [149], a radiation source that had previously been mostly overlooked. The contribution of environmental γ -rays to quasiparticle bursts was confirmed shortly afterwards in an experiment in which we exposed a transmon qubit to γ radiation sources [171]. By continuously monitoring the decay time of the qubit, that experiment was also the first to successfully operate a qubit as a γ radiation detector. Meanwhile, a second study on a fluxonium qubit found that these events, though detrimental to quantum computations, play a minor role in determining the average T_1 of qubits due to their rare occurrence. The development of a new firmware for our room-temperature electronics, and of a reset protocol that can work with any type of superconducting qubit, paved the way for the study of radiation-induced correlated errors in multi-qubit chips, one of the most worrisome phenomena in quantum computing. First tests with an LED illuminating a chip equipped with transmon qubits showed promising results.

The implications of this work, however, go beyond quantum computing. The experiment with the transmon proved that superconducting qubits represent a promising technology for the development of a new phonon-mediated particle detector, a possibility that had previously only been hypothesized. With an energy gap of $O(\text{meV})$ or lower, superconducting qubits could be sensitive to energy deposits too small for most of the particle detectors available today, making them an interesting candidate for future particle physics experiments.

Appendix A

A muon tagging system for superconducting quantum chips

In Chapter 3 it was illustrated how shielding superconducting quantum processors against cosmic muons is complicated because of their long range even in dense materials, with the most energetic ones penetrating kilometers of rock before being stopped. A straightforward solution to this problem would be moving the experiments to deep-underground facilities, where the rate of muon interactions is damped down to negligible values. This approach, however, has the drawback of requiring significant logistical and infrastructural effort to operate in such facilities.

An alternative is to implement in the experiment a tagging system that detects muons impacting the chip and discards any computation performed immediately afterward. The easiest realization of this system would consist of placing particle detectors outside the dilution refrigerator, but it would result in a trigger rate of hundreds of Hz [161], most of which would not compromise qubit operations at all (quantum processors have surfaces of a few cm^2 at most, meaning that the majority of the muons crossing the cryostat would not hit the chip inside). A much more efficient strategy, instead, consists of placing cryogenic detectors in the vicinity of the chip, which would then trigger only on muons that actually disturb the qubits.

Following this second strategy, in February 2025 we tested a demonstrator of a muon tagging system based on Kinetic Inductance Detectors (KIDs). The system consisted of two KIDs, each deposited on $4.5 \times 4.5 \text{ cm}^2$, $525\text{-}\mu\text{m}$ -thick silicon substrates that acted as target material for muon interactions. The two detectors were designed to be placed respectively above and below the quantum processor, following the idea that a muon crossing both of them would have also crossed the processor, compromising its operation. To better estimate the efficiency of this setup, however, we decided to replace

the quantum processor with a third KID detector, with significantly higher detection efficiency for particle interactions. From the analysis of the data, we concluded that our system achieves a muon-tagging efficiency of 90% with negligible dead time, confirming its potential.

Data acquisition, analysis and results have been summarized in an article titled *A Cryogenic Muon Tagging System Based on Kinetic Inductance Detectors for Superconducting Quantum Processors* [138], published on *Quantum Science and Technology*. The full text of the article is reported below.

A.1 Introduction

Superconducting quantum processors are among the most promising platforms for scalable quantum computation, offering design flexibility, high-fidelity operations, and compatibility with fast readout architectures [23, 32]. Significant advancements in the past decade include the demonstration of quantum advantage [17, 125, 185, 186], rapid progress in scaling up the number of qubits [187–190], and substantial improvements in qubit coherence times [89, 90, 108, 191]. These developments are key milestones toward fault-tolerant quantum computing.

However, ionizing radiation has recently been recognized as a source of decoherence in superconducting qubits, limiting device performance [118]. Energy deposited by environmental γ -rays or cosmic-ray muons in the chip substrate generates free charges, whose recombination produces high-energy phonons. These phonons propagate across the device, breaking Cooper pairs and producing bursts of quasiparticles. Quasiparticle tunneling across the Josephson Junction can significantly degrade the relaxation time of single qubits and, in some cases, induce correlated errors across multiple qubits on the same chip [131, 133, 136, 137]. These errors violate the assumptions of spatial and temporal independence that underlie most quantum error correction algorithms and represent a major challenge to large-scale quantum computing. Multiple studies have demonstrated a direct correlation between radiation exposure and qubit performance degradation, with event rates consistent with ambient radiation levels [118, 171]. Other investigations, however, suggest that ionizing radiation may also couple to two-level systems (TLSs) [139], a known dominant noise source in superconducting circuits.

To mitigate these effects, a variety of strategies have been explored. Careful material selection, passive shielding, and the use of underground laboratories can reduce radiation-induced instabilities and improve superconducting device performance [119, 135, 192].

However, deploying superconducting processors in deep-underground facilities introduces substantial infrastructure complexity. Other promising approaches include phonon “traps” [142, 143, 193], gap engineering [145, 194] to reduce quasiparticles tunneling, and even the integration of external radiation sensors [195]. While the interaction rate in a superconducting quantum chip is dominated by γ -rays from environmental radioactivity, which outnumber atmospheric muons by at least a factor of two [149, 196, 197], muons are particularly problematic due to their higher energy and penetration power. As a result, they cannot be attenuated by passive shielding and represent a dominant source of irreducible correlated errors in quantum devices operated above-ground.

In this work, we explore an alternative strategy based on a cryogenic muon-tagging system that can be seamlessly integrated into existing superconducting quantum processors. The system is compatible with standard cryogenic wiring and multiplexed RF readout schemes, and can share the same data acquisition electronics typically used for qubit operation. This intrinsic compatibility allows straightforward deployment within current cryogenic setups without requiring any modification to the existing hardware. We present the design, simulation, and experimental validation of the first prototype of such a muon-tagging system based on Kinetic Inductance Detectors (KIDs). The purpose of the present work is to demonstrate the particle-tagging capability of the system in a fully cryogenic environment, which represents a key prerequisite for its integration with sensitive quantum devices. The device operates at millikelvin temperatures and identifies muon events through coincident signals in a pair of KIDs positioned above and below the chip, respectively. We demonstrate efficient tagging of atmospheric muons ($\sim 90\%$) with negligible dead time, and show that the measured coincidence rate is in excellent agreement with Monte Carlo predictions. These results pave the way for the deployment of active veto architectures in above-ground quantum computing platforms, enabling the identification and, potentially, the rejection or correction of muon-induced correlated errors in real time.

A.2 System Design and Implementation

The muon-tagging system is designed to identify cosmic-ray muons interacting with a superconducting quantum chip by detecting coincident signals in adjacent superconducting sensors. These signals are produced by the energy deposited in the detector substrates as the particles traverse them. The system must be compatible with cryogenic operation, minimize electromagnetic interference, and ensure both mechanical and thermal

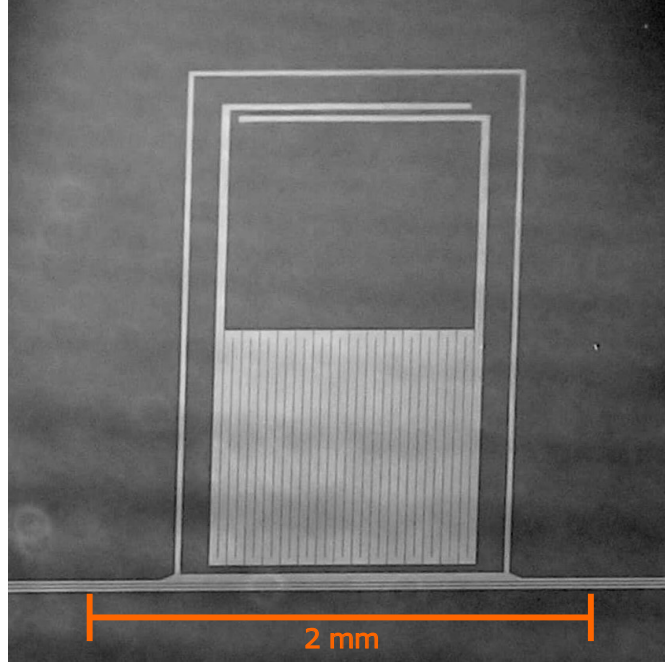


FIGURE A.1: Microscope image of a kinetic inductance detector (KID). The device consists of a meandered superconducting inductor (~ 6 cm in total length and $62.5 \mu\text{m}$ in width) connected to a two-finger interdigitated capacitor.

integration with the quantum device. In addition, it must cover a sufficiently large surface area to achieve high tagging efficiency while keeping the accidental coincidence rate negligible [198].

The experimental setup consists of a vertical stack of three detectors. The top and bottom layers form the muon-tagging system, while the central one acts as a proxy for the qubit chip during prototype testing. Each detector includes a $525\text{-}\mu\text{m}$ -thick, high-resistivity ($> 10 \text{ k}\Omega\cdot\text{cm}$) silicon substrate equipped with a KID. Each KID is a superconducting resonator consisting of a meandered inductor, approximately 6 cm in total length and $62.5 \mu\text{m}$ in width, connected to an interdigitated capacitor with two fingers, whose length sets the resonant frequency (Figure A.1). All devices were fabricated by the nanofabrication team at the Superconducting Quantum Materials and Systems (SQMS) Center at Fermilab (USA). Prior to metal deposition, the native silicon oxide was removed using a buffered HF dip. A 60 nm-thick aluminum film was then deposited on a 4-inch silicon wafer using electron-beam evaporation. The resonators, inductively coupled to a coplanar waveguide for readout, were defined by photolithography and etched into the aluminum layer using a chlorine-based dry etch. After etching, the wafer underwent standard photoresist stripping and cleaning before being diced with a precision saw according to the design layout. Each chip was then individually cleaned and packaged for integration into the experimental setup.

When a muon (or other particle) deposits energy in the silicon substrate, athermal phonons are generated and absorbed in the superconducting KID film, breaking Cooper pairs into dissipative quasiparticles. This changes the kinetic inductance, producing a shift in the resonant frequency proportional to the deposited energy [199, 200]. Thanks to the large absorber surface of the silicon wafer (20.25 cm^2) compared to the active area of the KID ($\sim 3.8 \text{ mm}^2$), the phonon-mediated detection approach provides high geometrical efficiency for muon tagging, achieving $\sim 90\%$ efficiency with only two devices (Sec. A.3). By contrast, hundreds of KIDs would be required to achieve comparable coverage around a qubit chip through a typical direct-absorption approach, where the incident particle deposits energy directly into the superconducting resonator [201, 202]. Although the phonon-mediated approach suffers from limited phonon collection efficiency, measured to be only a few percent for aluminum KIDs with this geometry [203], this limitation is not critical due to the high sensitivity of KIDs, which have energy thresholds of the order of 1 keV, well below the average energy deposited by muons in a $525\text{-}\mu\text{m}$ -thick silicon wafer ($\approx 150 \text{ keV}$). Phonon-mediated detection with KIDs is therefore a promising technique and is under active development for particle detection applications [199, 200, 204].

Finally, the detectors are mounted in copper holders using Teflon supports that clamp the wafer along its edges. This configuration minimizes both mechanical stress and contact with the active area, which could otherwise introduce additional phonon-loss pathways. The three devices are arranged in a top–center–bottom configuration with approximately 4.5 mm spacing between each layer. This spacing provides an effective compromise between mechanical constraints and detector proximity: it is small compared to the chip dimensions, ensuring high muon-tagging efficiency while still allowing sufficient clearance for mounting and wiring. The top and bottom detectors, forming the muon-tagging system, have a square geometry of $4.5 \times 4.5 \text{ cm}^2$, while the central layer hosts a KID on a $4 \times 2 \text{ cm}^2$ substrate. The central detector was intentionally designed with this geometry to emulate the footprint of a superconducting quantum processor. This area ($\sim 8 \text{ cm}^2$) is representative of a chip hosting on the order of 80–100 qubits, although the exact dimensions depend on the qubit layout and specific processor geometry. This choice allows testing the muon-tagging geometry under conditions representative of a real quantum device. In our applications, which employ smaller chips (typically $< 1 \text{ cm}^2$), an even higher tagging efficiency is expected, since the external detectors would cover a larger fraction of the active area. Figure A.2 shows the three detectors and their assembly in the vertical stack.

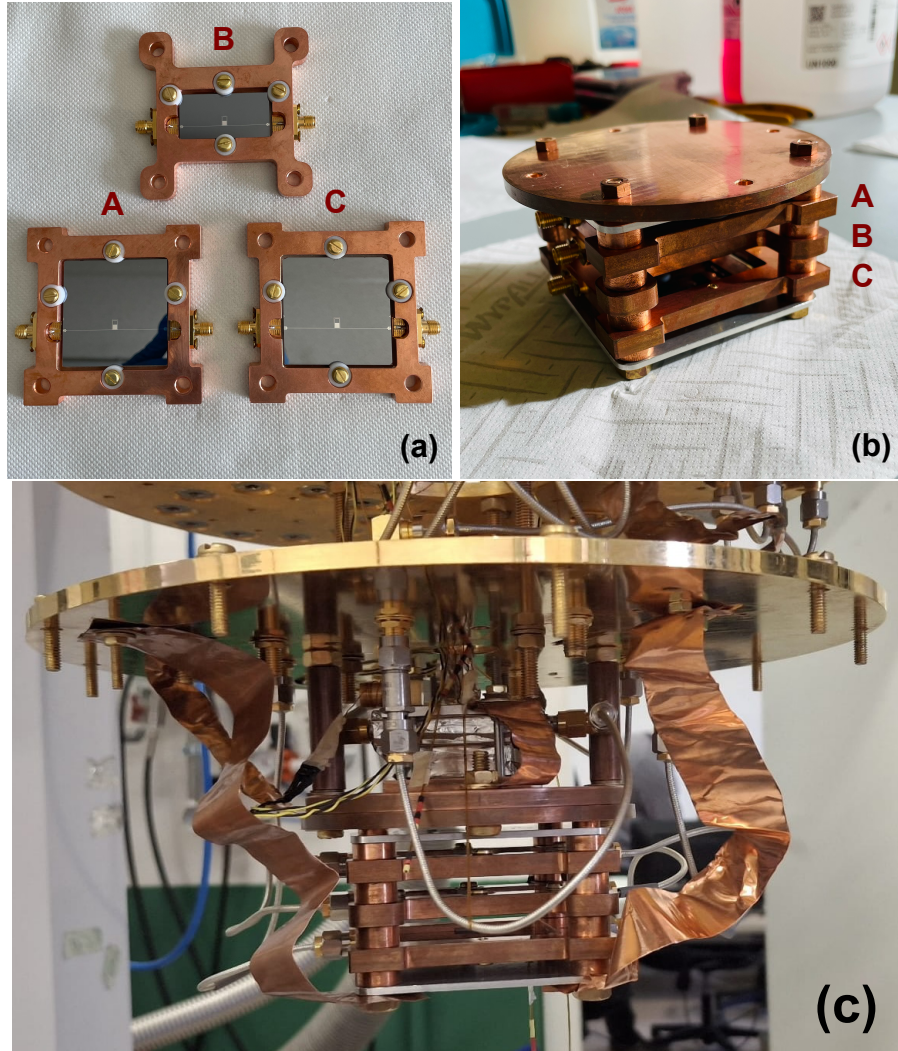


FIGURE A.2: (a) Picture of the three KID-based detectors: the top (A) and bottom (C) detectors form the muon-tagging system, while the central detector (B) serves as a proxy for the qubit chip in the prototype. (b) Vertical stack of the detectors mounted in the copper holder, showing the top-center-bottom configuration with a 4.5 mm inter-layer spacing. (c) The assembled detector stack installed at the mixing-chamber stage of the dilution refrigerator at La Sapienza University.

A.3 Monte Carlo Simulations

To evaluate the expected performance of the muon-tagging system and to guide the prototype design, a Monte Carlo simulation was developed using the Geant4 toolkit [146]. The simulation implements the full three-dimensional geometry of the detector stack, the copper holders, and a simplified model of the cryostat including a three-layer shielding pot (Al/Cu/Cryophy®) described in Sec. A.4. The three devices are modeled as 525- μm -thick silicon substrates with the same dimensions as the fabricated detectors, arranged in the top-center-bottom configuration illustrated in Figure A.2c. In the model, the central layer, acting as a proxy for the qubit chip, is enclosed in a 5 mm-thick copper

box, while the top and bottom detectors are mounted in open copper holders, reproducing the experimental setup. This configuration replicates the mechanical mounting commonly adopted for superconducting quantum processors, where the chip is enclosed in a copper box for mechanical support and thermalization. In real qubit setups, the copper box is also surrounded by a thin mu-metal shield. Since such shielding affects only the magnetic environment and not the passage of minimum-ionizing particles, it is not included in the simulation and has no impact on the expected tagging performance.

The simulation tracks the passage of particles and their energy deposition in the silicon substrates. The subsequent phonon propagation and detector response are not modeled; instead, the deposited energy is used as a proxy for the production of detectable signals. Because KIDs are sensitive to any energy deposition in the substrate, signals produced by environmental γ -rays, electrons, or α particles are indistinguishable from those induced by muons. At high single-detector rates, this can lead to accidental coincidences within the coincidence window, increasing the system dead time and thus reducing the effective muon-tagging efficiency. Among these backgrounds, γ -rays dominate due to their higher flux compared to other particles [149], and were therefore simulated explicitly.

Primary particles, i.e., the ionizing particles injected into the simulation, were generated within the laboratory environment. Atmospheric muons were simulated to evaluate the tagging efficiency, while environmental γ -rays were included to estimate the γ -induced dead time. The muon spectrum and angular distribution reproduce measurements of cosmic-ray muons at sea level [205], whereas the γ flux was derived from measurements performed with a 3-inch portable NaI(Tl) scintillator in the laboratory where the prototype was tested (“Laboratorio di Rivelatori Criogenici” at La Sapienza University, Italy). Muons are generated from random positions on a hemispherical surface above the cryostat according to the method described in Ref. [149], while photons from ambient radioactivity are generated isotropically following the measured γ -ray energy distribution. Particles are then propagated through the full geometry, and the energy deposited in each silicon wafer is recorded. Interaction rates in the detectors are obtained by scaling the number of recorded events to the measured fluxes: $1 \mu/\text{cm}^2/\text{s}$ at sea level and $13.5 \gamma/\text{cm}^2/\text{s}$ in the laboratory, the latter being substantially higher than the typical $2\text{--}3 \gamma/\text{cm}^2/\text{s}$ observed in standard above-ground environments.

Table A.1 summarizes the simulated interaction rates for muons and γ -rays in the top and bottom detectors, along with the expected coincidence rates. For atmospheric muons, requiring simultaneous energy depositions in the two detectors yields a coincidence rate of $(195 \pm 12) \times 10^{-3}$ events/s. For environmental γ -rays, the coincidence rate is much lower than for muons, as their interaction probability in silicon is small and the chance of producing energy depositions in both detectors is significantly reduced. Figure A.3

Channel	Muon Rate [events/s]	Gamma Rate [events/s]	Total Rate [events/s]
T	$(312 \pm 19) \times 10^{-3}$	2.9 ± 0.2	3.2 ± 0.2
B	$(296 \pm 17) \times 10^{-3}$	3.0 ± 0.2	3.3 ± 0.2
T-B	$(195 \pm 12) \times 10^{-3}$	$(15 \pm 4) \times 10^{-3}$	$(210 \pm 13) \times 10^{-3}$

TABLE A.1: Simulated interaction rates for muons and γ -rays in the top (T) and bottom (B) detectors, and expected coincidence rates (T-B). Rates are normalized to fluxes of $1 \mu/\text{cm}^2/\text{s}$ and $13.5 \gamma/\text{cm}^2/\text{s}$, respectively. Uncertainties are dominated by a 5% systematic component from the Geant4 simulation.

shows the simulated energy deposition of atmospheric muons in the central chip and the subset of events tagged by the external detectors (top and bottom) forming the muon-tagging system. The ratio of tagged to total events ($4589/5073$) corresponds to a muon-tagging efficiency of about 90%.

In a coupled qubit–muon-tagging system architecture, every coincident event would activate a veto window during which quantum operations are rejected to avoid correlated errors due to muon interactions within the chip. The characteristic recovery time observed in superconducting qubits ranges from a few milliseconds [137, 171] up to 25 ms in some experiments [133], depending on the qubit geometry. It is therefore crucial to assess whether random coincidences between independent detectors could introduce a significant fractional dead time.

Using the rates in Table A.1 and noting that the contribution from μ – μ accidental coincidences is negligible, the accidental coincidence rate between the top and bottom detectors within a $340 \mu\text{s}$ coincidence window (corresponding to the selection applied in the data analysis) can be estimated as the sum of:

1. $R_{\text{acc}}^{\gamma\gamma} = R_{\text{T}}^{\gamma} R_{\text{B}}^{\gamma} \Delta t \approx (3.0 \pm 0.2) \times 10^{-3} \text{ events/s};$
2. $R_{\text{acc}}^{\gamma\mu} = (R_{\text{T}}^{\gamma} R_{\text{B}}^{\mu} + R_{\text{T}}^{\mu} R_{\text{B}}^{\gamma}) \Delta t \approx (0.6 \pm 0.1) \times 10^{-3} \text{ events/s};$

The total accidental-coincidence rate is therefore

$$R_{\text{acc}} \approx (3.6 \pm 0.2) \times 10^{-3} \text{ events/s.} \quad (\text{A.1})$$

The simulation also predicts a true γ – γ coincidence rate of $(15 \pm 4) \times 10^{-3} \text{ events/s}$ (see Table A.1). Assuming a veto gate of 1–5 ms, the resulting fractional dead time due to γ -induced coincidences is

$$f_{\text{DT}}^{(\gamma)} = (R_{\text{TB}}^{\gamma\gamma} + R_{\text{acc}}) \tau_{\text{veto}} \approx (2\text{--}9) \times 10^{-5}, \quad (\text{A.2})$$

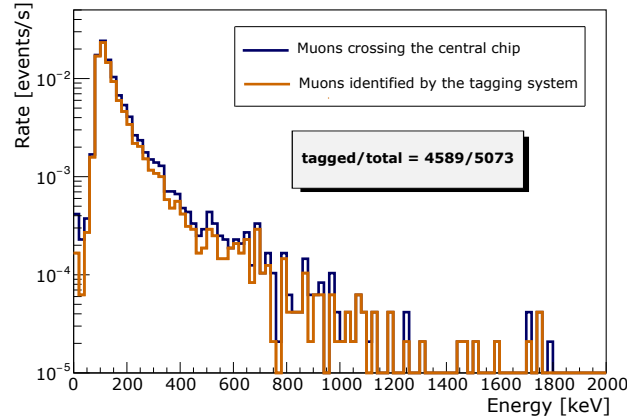


FIGURE A.3: Simulated energy deposition of atmospheric muons in the central chip (blue) and events identified by the muon-tagging system (orange). The ratio of tagged to total events (4589/5073) corresponds to a muon-tagging efficiency of about 90%.

i.e., well below 0.01%. Even in the extreme case of a 25 ms veto gate, the expected fractional dead time is 0.05%. Therefore, the contribution of γ -induced coincidences to dead time is entirely negligible, confirming that ambient γ -rays do not limit the live time performance of the system, even under high-flux conditions ($13.5 \gamma/\text{cm}^2/\text{s}$). This makes the present estimate a conservative upper bound for the expected dead time.

These results confirm that, despite the relatively large substrate area and the resulting higher single-detector event rate, the coincidence-based muon-tagging approach is not limited by accidental or true coincidences induced by environmental γ -rays. The simulation therefore validates the feasibility of operating such a system without significant live time losses, providing the baseline for the experimental results presented in the next sections.

A.4 Experimental Setup and Data Acquisition

The prototype was operated in a dry $^3\text{He}/^4\text{He}$ dilution refrigerator at a base temperature of approximately 20 mK at La Sapienza University. The vertical stack of the three detectors was enclosed in a three-layer shielding pot made of aluminum (inner), copper (middle), and Cryophy® (outer), which provided thermal shielding and protection from external magnetic fields. The assembly was thermally anchored to the mixing-chamber stage through copper bars for thermalization (Figure A.2c).

RF connectors and coaxial lines allowed the simultaneous readout of all three detectors in a daisy-chain configuration. An RF tone generated at room temperature by an Ettus X310 board was used to bias the KIDs, and attenuated by 55 dB along the cryostat to suppress thermal noise before reaching the array. The transmitted signal was amplified

by a HEMT low-noise amplifier at the 4 K stage and then returned to the same board for down-conversion. Data were acquired using the open-source firmware developed for the BULLKID experiment [206]. When a signal exceeded a predefined threshold, 24 millisecond-long waveforms from all channels were recorded simultaneously, with any detector able to act as trigger. The trigger condition was defined as a resonator response, after band-pass filtering, exceeding five times the RMS of the filtered baseline noise. Data were acquired with a sampling frequency of 100 kHz and a total live time of about one hour per run.

When integrated with a superconducting qubit chip, the KID-based muon tagging system would operate in parallel with the qubit control and readout electronics. While the KIDs are read out here using an Ettus X310 board, the qubits could be controlled and measured using an RFSoc-based platform such as a ZCU216 board [171]. The two readout chains would initially remain independent within the same cryogenic environment. Preliminary tests with a qubit chip featuring $T_1 \sim 100 \mu\text{s}$ demonstrate that potential electromagnetic crosstalk between the muon-tagging system and the qubits does not introduce any observable degradation of qubit T_1 . Synchronization between the KID and qubit acquisition systems can be achieved via optical calibration pulses injected through a fiber-coupled LED down to the mixing chamber, enabling temporal alignment of the datasets. As a first step, the veto could be applied offline by discarding qubit data within a defined time window following a KID trigger. Radiation-induced events in similar devices typically last 1–2 ms [171], providing a guideline for defining the veto window.

By fitting all the resonators using the method described in Ref. [207], we found resonance frequencies f_0 consistent with the design, in the 2.52–2.57 GHz range, with coupling quality factors $Q_c \sim 10^5$ and internal quality factors $Q_i \sim 10^5 - 10^6$, depending on the device. When an amount of energy E is released into the silicon wafer, changes in the Cooper-pair density are detected by monitoring the time variation of S_{21} with the RF bias set at $f = f_0$. Although both magnitude and phase of S_{21} can be used, the phase typically yields a better signal-to-noise ratio.

Pulses were reconstructed with a matched filter optimized for each channel [208], and only those exceeding five times the baseline RMS in each channel were retained to suppress spurious triggers. More refined strategies for setting the energy threshold have also been proposed [209], but are beyond the scope of the present study. Particle-induced pulses exhibited decay times of several hundred microseconds, enabling effective separation of single events and sub-millisecond time resolution for coincidence analysis. Coincidences were identified by selecting a pulse maximum in one detector (reference channel) and verifying that the maximum in the other detector occurred within a $\pm 170 \mu\text{s}$

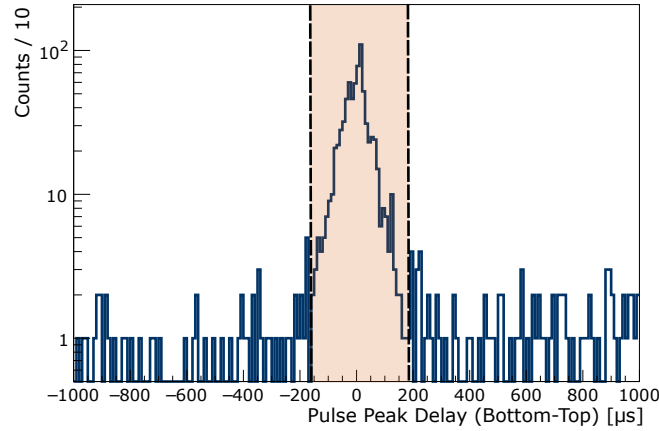


FIGURE A.4: Distribution of the *Pulse Peak Delay* between the bottom and top detectors. The shaded area ($|\Delta t| \leq 170 \mu\text{s}$) represents the $340 \mu\text{s}$ -wide coincidence window used to select true coincident events, while the flat regions at larger delays correspond to random coincidences (sidebands) employed to estimate the accidental background.

time window, resulting in a total coincidence window of $340 \mu\text{s}$. The distribution of time differences between the pulse maxima, defined as Pulse Peak Delay (Bottom–Top), is shown in Figure A.4, where the central shaded region marks the chosen coincidence window. The $\pm 170 \mu\text{s}$ limit corresponds to three standard deviations of the Gaussian fit to the coincidence-time peak in the data and is also consistent with the characteristic rise time of the KID pulses. Due to the finite rise time of the detectors, the temporal ordering of energy depositions in the top and bottom devices cannot be reliably inferred from the pulse maxima. The coincidence window is therefore defined to identify events occurring in both detectors, without attempting to determine their sequence. Figure A.5 shows the pulse amplitudes in the two detectors before and after applying the amplitude and time-selection cuts. In the pre-selection sample, single-detector triggers and uncorrelated events populate the region at low amplitude in one of the two detectors. After the cuts, a cluster of true coincidences becomes clearly visible. Figure A.6 displays an example of a muon event producing simultaneous signals in all three detectors. The three detectors exhibit different pulse amplitudes and baseline noise levels, as expected for independent KID devices. While individual pulse amplitudes fluctuate on an event-by-event basis, the baseline RMS remains stable throughout the data-taking period, ensuring reliable operation without the need for frequent recalibration. Selection thresholds are defined in units of each detector’s RMS, providing a consistent signal-to-noise requirement across the system and avoiding bias in coincidence identification. The relative performance of the detectors is broadly consistent across the dataset, supporting the robustness of the coincidence selection.

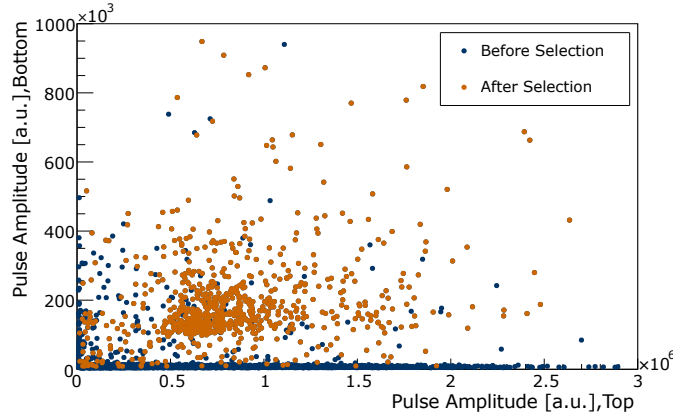


FIGURE A.5: Scatter plot of pulse amplitudes in the top and bottom detectors before (blue) and after (orange) applying the selection cuts described in the text. Before the cuts, the dataset is populated by noise, single-detector triggers, and uncorrelated events, which appear as events with negligible amplitude in one of the two detectors. After the cuts, a cluster of true coincidences stands out more clearly.

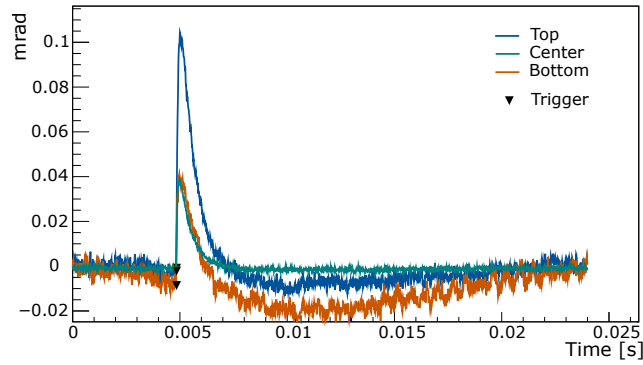


FIGURE A.6: Example of a muon event producing coincident signals in the three detectors.

A.5 Results

A.5.1 Single-detector rates

Applying the baseline RMS selection cut described in Sec. A.4 to the reconstructed pulse amplitudes yields the single-detector rates reported in Table A.2. The measured values are in good agreement, within uncertainties, with the simulated interaction rates reported in Table A.1.

A.5.2 Accidental coincidences

Using the measured single-detector rates and a $340\,\mu\text{s}$ coincidence window, the expected rate of random coincidences between the top and bottom detectors is $R_{\text{acc}}^{\text{exp}} =$

Channel	Measured Rate [events/s]	Expected Rate [events/s]
T	3.43 ± 0.03	3.2 ± 0.2
B	2.97 ± 0.03	3.3 ± 0.2
T-B	$(211 \pm 8) \times 10^{-3}$	$(214 \pm 13) \times 10^{-3}$

TABLE A.2: Measured single-detector and coincidence rates after applying the baseline RMS selection cut, compared with the simulated interaction rates from Table A.1. The T-B coincidence rate includes contributions from muon, γ - γ , and accidental coincidences. Measurement errors are statistical only, while simulated rates are dominated by a 5% systematic uncertainty from the Geant4 simulation.

$R_T^{\text{meas}} R_B^{\text{meas}} \Delta t \approx (3.5 \pm 0.1) \times 10^{-3}$ events/s, in excellent agreement with the Monte Carlo prediction $((3.6 \pm 0.2) \times 10^{-3}$ events/s).

The rate of random coincidences was also evaluated using the flat sidebands in Figure A.4, corresponding to time delays outside the coincidence region. Because the acquisition window is asymmetric—with a longer post-trigger region (see Figure A.6)—only the negative-delay sideband was used to avoid bias. Integrating events in the interval $-2500 \leq \Delta t \leq -400 \mu\text{s}$ to reduce statistical fluctuations, and scaling to the $340 \mu\text{s}$ coincidence window, yields $R_{\text{acc}}^{\text{SB}} = (3.9 \pm 0.4) \times 10^{-3}$ events/s, fully consistent with both the analytical and simulated expectations.

This analysis confirms that the accidental coincidence rate is negligible compared to the total coincidence rate, ensuring that the observed excess in the coincidence window originates from true muon events.

A.5.3 Muon-induced coincidences

Applying the coincidence requirement described in Sec. A.4 yields a total top-bottom coincidence rate of $R_{\text{T-B}}^{\text{tot}} = (211 \pm 8) \times 10^{-3}$ events/s, as reported in Table A.2. Subtracting the expected contribution from γ -induced and accidental coincidences ($R_{\text{T-B}}^{\gamma\gamma} + R_{\text{acc}} = (18.6 \pm 4.0) \times 10^{-3}$ events/s) gives a muon-induced coincidence rate of

$$R_{\text{T-B}}^{\mu\mu} = (192 \pm 9) \times 10^{-3} \text{ events/s},$$

in excellent agreement with the Monte Carlo prediction of

$$R_{\text{T-B}}^{\mu\mu, \text{MC}} = (195 \pm 12) \times 10^{-3} \text{ events/s}.$$

This correspondence confirms that the system achieves the expected muon-tagging efficiency of approximately 90%. Given the measured single-detector rates and coincidence window, the resulting fractional dead time due to γ -induced coincidences remains below

10^{-4} , demonstrating that the tagging performance is unaffected by background events and that the system can operate with negligible live time losses.

A.6 Conclusions and Outlook

We have presented the design, implementation, and first operation of a cryogenic muon-tagging system based on kinetic inductance detectors, developed as a proof-of-principle platform for future integration with superconducting quantum processors. The system consists of a vertical stack of two KIDs deposited on large area substrates, operated at millikelvin temperatures and readout through a multiplexed RF architecture adapted from the BULLKID experiment. A third, smaller detector placed in the middle acts as a proxy for the qubit chip during prototype testing.

Monte Carlo simulations performed with Geant4 guided the system design and established quantitative expectations for the muon-tagging efficiency as well as for both true and accidental coincidences induced by ambient γ -rays. In practical implementations for research purposes, the central module would host a superconducting qubit chip typically with dimensions of the order of 1 cm^2 . Under such conditions, an even higher tagging efficiency would be achieved, since the external detectors would cover a larger fraction of the active area. In our prototype, instead, the intermediate KID was intentionally fabricated with a significantly larger area (8 cm^2), comparable to that of a multi-tens-qubit processor, to emulate a conservative upper-bound scenario. Conversely, larger quantum processors could still maintain efficiencies above 90% by extending the muon-tagging system coverage, with only a marginal increase in dead time thanks to the low background rate demonstrated here.

The fabricated prototype achieved internal quality factors up to 10^6 and pulse decay times of several hundred microseconds, enabling sub-millisecond time resolution for coincidence analysis. Using the selection criteria described in Sec. A.4, we measured a top-bottom muon coincidence rate of $(192 \pm 9) \times 10^{-3}\text{ events/s}$, in excellent agreement with the Monte Carlo prediction. The corresponding muon-tagging efficiency is about 90%, with a fractional dead time below 10^{-4} , confirming that ambient γ -rays do not limit the live time performance of the system.

These results demonstrate the feasibility of operating a high-efficiency, low-dead-time muon-tagging system at millikelvin temperatures, establishing the key technological prerequisite for its integration into above-ground superconducting quantum processors. Future developments will focus on coupling the tagging system to a multi-qubit chip, introducing veto strategies or correction algorithms against muon-induced correlated errors,

and optimizing KID geometries and materials (e.g., inductor volume, capacitor design, feedline coupling) to further improve detection efficiency.

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