Rapid Driven Reset of a Qubit Readout Resonator

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Using a circuit QED device, we demonstrate a simple qubit measurement pulse shape that yields fast ring-up and ring-down of the readout resonator regardless of the qubit state. The pulse differs from a square pulse only by the inclusion of additional constant-amplitude segments designed to effect a rapid transition from one steady-state population to another. Using a Ramsey experiment performed shortly after the measurement pulse to quantify the residual population, we find that compared to a square pulse followed by a delay, this pulse shape reduces the timescale for cavity ring-down by more than twice the cavity time constant. At low drive powers, this performance is achieved using pulse parameters calculated from a linear cavity model; at higher powers, empirical optimization of the pulse parameters leads to similar performance.

Over the last decade, circuit quantum electrodynamics (cQED) has become the leading architecture for constructing scalable networks of solid-state qubits, finding application in the context of not only superconducting qubits [2–3] but also spin qubits [4] and potentially others [5]. In this paradigm, each qubit is coupled to a resonator in which it induces a state-dependent frequency shift, allowing the qubit state to be interrogated using a pulsed tone near the resonator frequency. Although improvements in quantum-limited amplifiers [6–9] have made fast (sub-microsecond), high-fidelity readout in cQED systems routine, relatively little attention has been devoted to the problem of rapidly returning the readout resonator to its ground state immediately after the measurement pulse. If the measurement pulse is simply turned off, residual photons gradually exiting the resonator will continue to measure the qubit [10], preventing high-fidelity operations. Although the time constant for this decay can be shortened by using a resonator with a broad linewidth, doing so increases the rates of both decay via spontaneous emission (Purcell effect [11]) and dephasing from thermal photons [12, 13]. While the former effect may be mitigated using a Purcell filter [14, 15], the latter may only be compensated by reducing the qubit-resonator coupling or the resonator temperature. A technique for reducing the residual population on a timescale faster than the resonator’s free decay will therefore be important in any algorithm in which re-using qubits shortly after measurement is important, e.g., error correction with surface [16, 17], C4 [18] or Bacon-Shor [19] codes. A major impediment has been the fact any such scheme typically derived CLEAR pulse shape depopulates the cavity to a negligible level in a time more than two cavity time constants faster than that needed after the square pulse. At higher powers, optimizing the pulse shape empirically using an iterative algorithm leads to equally good performance.

The experimental device is a fixed-frequency transmon qubit mounted in a 3D aluminum cavity [20] attached to the mixing chamber of a dilution refrigerator at an indicated base temperature of 10 mK. Qubit and measurement drive tones are generated using Agilent E8267D function generators and modulated using Tektronix AWG7000 series arbitrary waveform generators at a 2 GS/s sample rate. Qubit pulses are 4σ Gaussians with DRAG [21] correction. The cavity is measured in transmission, and the transmitted signal is fed to a HEMT amplifier (Low Noise Factory LNF-LNC620A) at 4 K using a superconducting NbTi/NbTi semi-rigid coaxial cable. After additional amplification at room temperature, the signal is mixed down to 16 MHz and digitally demodulated. The cavity is measured to have bare frequency $f_{\text{bare}} = 10.7457$ GHz, dressed frequency $f_{\text{dressed}} = 10.7588$ GHz, and linewidth $\kappa/2\pi = 1.1$ MHz (corresponding to a time constant $1/\kappa = 0.14$ μs). The qubit has frequency $f_{01} = 4.83315$ GHz, anharmonicity $\delta/2\pi = -151$ MHz, average $T_1 \approx 50$ μs, and average $T_2^{\text{echo}} \approx 60$ μs. Preparing the qubit in the excited state shifts $f_{\text{dressed}}$ by the cavity pull $2\chi/2\pi = -2.6$ MHz; the measurement tone is applied at the midpoint of the two frequencies, $f_{\text{dressed}} - \chi/2\pi$.

The residual population after a measurement pulse can be quantified in terms of the mean cavity photon number $n$ at some time after the end of the pulse. We use the
sequence illustrated in Fig. 1(b) to extract \( n \) following an initial measurement pulse denoted M1. The measurement pulse is followed by a quick Ramsey experiment (\( t_{\text{gate}} = 8 \text{ ns}, t_{\text{R}} = 0 \text{ to } 600 \text{ ns} \)) to probe the ac Stark shift and dephasing from any residual photons \[10][22]. The time \( t_{\text{relax}} \) between the end of M1 and the start of the Ramsey experiment can be varied to measure \( n \) as a function of time after the end of M1. The time \( t_{\text{buffer}} \) between the Ramsey experiment and the second measurement pulse (M2) is set to 400 ns to ensure that even when \( t_{\text{g}} \) and \( t_{\text{relax}} \) are both short, M2 is not corrupted by lingering photons from M1.

A typical Ramsey trace is shown in Fig. 2(a). The non-monotonic modulation in both amplitude and frequency arises from the fact that the cavity population evolves during the Ramsey delay, leading to recurrences. We derive the expected form of this transient response using the positive-P-function method as in Gambetta et al. \[10\], where it was applied to the steady-state problem. For a Ramsey detuning \( \Delta \) (here 10 MHz), decoherence occurs when \( \Gamma \) becomes large, and the only free parameters are \( n_0 \) and \( \phi_0 \). As illustrated in the figure, this function yields a good fit to the data, allowing reliable determination of \( n_0 \). In the rest of this work, we use the extracted \( n_0 \) to quantify the residual population as a function of pulse shape, drive power, and wait time. We note that \( n_0 \) does not include the background thermal population of the cavity, which is accounted for by the \( T_2^\text{echo} \) term and is calculated from \( T_2^\text{echo} \) to be \( \sim 0.02 \) on average.

The Ramsey fit method of obtaining \( n_0 \) was validated by using it to measure \( n_0 \) as a function of both wait time \( t_{\text{relax}} \) and pulse power. For simplicity, these tests were performed using square pulses for both M1 and M2, varying only the power of M1, denoted \( P \) as illustrated in Fig. 1(a). For convenience we define the normalized drive power as \( P_{\text{norm}} = P/P_{\text{1ph}} \), where \( P_{\text{1ph}} \) is the steady-state drive power that yields \( n = 1 \), as inferred from a standard Ramsey experiment measuring the Stark shift \( \Delta \omega = 2 \chi n \[10\] \) induced by a CW tone. Figure 2(b) shows \( n_0 \) extracted from Ramsey fits as a function of \( t_{\text{relax}} \). Regardless of the prepared qubit state, the decay is exponential, as expected for free decay, and the time constant \( \tau \) extracted from the best-fit curve is consistent with the value of \( \kappa \) obtained from frequency-domain measurements. Figure 2(c) shows \( n_0 \) as a function of \( P_{\text{norm}} \).
Residual population $n_0$; similar behavior was observed at all values of $t_{relax}$ for which non-negligible $n_0$ were measured. The dashed line indicates the expected behavior assuming a linear cavity: $n_0 = e^{-\kappa t_{relax}}P_{\text{norm}}$. The data exhibit a transition from a linear response at low powers to a super- (sub-) linear response at high powers when the qubit is prepared in the ground (excited) state. This behavior is consistent with the cavity’s expected self-Kerr nonlinearity [23], which shifts the cavity frequency by a negative amount $K$ per cavity photon, pushing it closer to (farther from) the measurement frequency when the qubit is in the ground (excited) state. Approximating $K$ in the small-$\delta$ limit as $K = 2g^4\delta(3\omega_q^4 + 2\omega_q^2\omega_r^2 + 3\omega_r^4)/(\omega_q^2 - \omega_r^2)^4$, where $\omega_q = f_{01}/2\pi$, $\omega_r = f_{\text{dressed}}/2\pi$, and $g$ is calculated as in Ref. [12], we calculate $K \approx -65$ kHz. We expect the linear cavity model to break down as $n$ approaches $|\kappa/K|$, or approximately 15 for our parameters.

Having thus validated our method of quantifying residual photons, we then switched to a CLEAR pulse shape for M1; for consistency, we continued to use a square pulse for M2. For an ideal single qubit-cavity system, the optimal CLEAR pulse envelope [Fig. 3(a)] consists of five piecewise-constant segments: two ring-up segments, one steady-state segment, and two ring-down segments. For the pulse to behave as intended, its bandwidth needs only to be much greater than that of the cavity, a condition readily achieved in our setup. Setting the carrier frequency to $f_{\text{dressed}} - \chi/2\pi$, as done here, is not required but maximizes both SNR and simplicity: in this case, a given measurement pulse yields the same steady-state $n$ regardless of qubit state as long as the cavity remains in the linear-response regime. The lengths of the ring-up and ring-down segments were all initially fixed at 150 ns (approximately $1/\kappa$), and their amplitudes relative to that of the steady-state segment were calculated by solving a driven damped harmonic oscillator model to find the pulse shape that populates and depopulates the cavity in the shortest amount of time regardless of the qubit state.

The cavity IQ plane trajectories produced by square and CLEAR pulses with the same steady-state amplitude ($P_{\text{norm}} = 3.6$) are shown by markers in Figs. 3(a,b), respectively. For all trajectories, the time step between markers is 24 ns. Solid lines in these plots indicate the theoretically calculated response of the cavity to each pulse, multiplied by an overall amplitude factor to match the data and adjusted to reflect the independently observed 20% thermal population of the qubit excited state (which was reduced by an order of magnitude on later cooldowns of this device with additional input line attenuation). We see that the experimental cavity responses track the theoretically calculated ones very well, and that compared to the square pulse, the CLEAR pulse yields more compact trajectories that reach near-steady-state populations (at both $n \approx 3.6$ and $n \approx 0$) in less time.

We quantitatively compare the performance of the CLEAR pulse to that of a square pulse using the Ramsey fit method. To provide a fair comparison, a zero-amplitude segment is appended to the square pulse to allow undriven decay during a time equivalent to the total length of the CLEAR pulse’s two ring-down segments. The results are shown in Fig. 3(c). At all measurement powers, the CLEAR pulse significantly outperforms the square pulse; moreover, for drive powers that keep the cavity in the linear regime (evidenced by $n_0$ being independent of the prepared qubit state), the residual population after the CLEAR pulse is negligible. At higher powers, it appears that the cavity nonlinearity, not taken into account in calculating the optimal CLEAR pulse parameters, prevents perfect ring-down. We also find non-idealities when the lengths of the ring-down segments are reduced in an effort to shorten the ring-down time: for 120 ns ring-down segments (a 20% reduction), we find measurable $n_0$ even in the linear regime, as illustrated.

FIG. 3. (color) (a,b) IQ-plane cavity trajectories in response to a square pulse and CLEAR pulse, respectively. In each plot, experimental results (markers) are superimposed on theoretical calculations (solid curves), the dashed circle indicates the target population $n = 3.6$, the black cross indicates the origin, and the black arrows indicate the directions of the trajectories. (c) Residual cavity population versus pulse power for both square and CLEAR pulse shapes. For the CLEAR pulse, the Ramsey experiment begins immediately at the end of the pulse, while for the square pulse, a delay of approximately 300 ns is inserted to match the total length of the CLEAR pulse’s two ring-down segments.
in Fig. 4(a). Further reductions in the segment lengths increase the performance degradation.

![Graphs](image)

**FIG. 4.** (color) (a) $n_0$ versus drive power for the shortened CLEAR pulse (120 ns ring-down segments), for $t_{\text{relax}} = 0$. (b) Evolution of $n_0$ at each step of an empirical optimization algorithm for the shortened CLEAR pulse with $P_{\text{norm}} = 10$. Inset: shortened CLEAR pulse shape before optimization (magenta) and after (green). (c) Ramsey traces obtained using initial shortened CLEAR pulse, yielding $n_0 \approx 2.2$ for ground and $n_0 \approx 0.91$ for excited. (d) Ramsey traces obtained using final shortened CLEAR pulse, yielding $n_0 < 0.1$ for both qubit states.

To improve performance of the CLEAR pulse both at high powers and with shortened ringdown segments, we use an empirical technique to optimize the pulse parameters. As an example, we take as a starting point the pulse with 120 ns ring-down segments and a steady-state drive power of $P_{\text{norm}} = 10$ (which yielded $n_0 \approx 2.16$ for the ground state and $n_0 \approx 0.91$ for the excited state), and run an iterative optimization algorithm that attempts to minimize $n_0$ by adjusting the amplitudes of the ring-down segments. We keep $t_{\text{relax}} = 0$ throughout. The evolution of $n_0$ with each iteration is shown in Fig. 4(b): in fewer than 300 iterations, the pulse is optimized to yield $n_0 < 0.1$ regardless of initial qubit state. Ramsey experiments before and after running this optimization are shown in Fig. 4(c,d), revealing that coherence is preserved with the optimized pulse shape regardless of initial qubit state. As seen in the inset of Fig. 4(b), the optimization process significantly increases the amplitudes of both ring-down segments. Extending our theoretical calculations to the non-linear regime may shed light on this result and potentially eliminate the need for empirical tune-up in this regime.

In summary, using a single-qubit cQED system, a qubit-state-independent reduction in the time needed to reach a steady-state resonator population both during and after a qubit measurement pulse was achieved by including extra constant-amplitude segments in the pulse. For low-power drives, near-perfect ring-down (quantified using Ramsey experiments) is achieved using segment amplitudes calculated from system parameters. At higher drive amplitudes, similar performance is obtained following empirical optimization of the pulse shape.

Though this demonstration used a 3D transmon, the same technique should be applicable to any cQED system; it may also be combined with machine-learning based analysis [24] and Purcell filters to further reduce the measurement cycle time. Future areas of interest may include numerical calculation of the optimal CLEAR parameters in the non-linear regime, extension of this technique to resonators coupled to multiple qubits, and investigation as a possible method for implementing the resonator-induced phase (RIP) gate [23] non-adiabatically.

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