Driven-state relaxation of a coupled qubit-defect system in spin-locking measurements

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It is widely known that spin-locking noise-spectroscopy is a powerful technique for the characterization of low-frequency noise mechanisms in superconducting qubits. Here we show that the relaxation rate of the driven spin-locking state of a qubit can be significantly affected by the presence of an off-resonant high-frequency two-level-system defect. Thus, both low- and high-frequency defects should be taken into account in the interpretation of spin-locking measurements and other types of driven-state noise-spectroscopy.

The field of gate-based quantum computing using superconducting qubits is rapidly developing [1]. However, the realization of a fault-tolerant quantum computer remains a challenging task: the implementation of a reasonably robust logical qubit would require an overhead of order $10^3$ to $10^4$ physical qubits with per-gate error rates $p \approx 10^{-3}$ [2]. Lowering per-operation error rates would significantly reduce the overhead requirements, and, therefore, it is important to understand noise-induced error mechanisms during free and driven evolution of qubits. Free-evolution qubit relaxation has been extensively studied, and standard protocols for measurements of energy relaxation and dephasing rates — including spin-echo, Ramsey, and dynamical decoupling methods — have been widely used [3, 4]. As for the decoherence of a superconducting qubit during driven evolution, qubit relaxation mechanisms in Rabi and spin-locking measurements were initially analyzed in Refs. [5, 6]. The Rabi spectroscopy was later used in experimental studies of noise characteristics in superconducting flux qubits [7, 8], and in a qubit-fluctuator system [9]. More recently, spin-locking spectroscopy has been demonstrated to be a powerful tool for noise characterization in superconducting qubits [10], which can be used to detect low-frequency defects [11], distinguish between coherent and thermal photon noise [12], and measure low-frequency noises in multi-qubit systems [13].

Spin-locking noise spectroscopy of a superconducting qubit is based on the measurements of the qubit evolution driven by a spin-locking pulse sequence originally developed in NMR studies [10]. In the Bloch-sphere representation in a rotating frame, a qubit state is associated with a fictitious spin bi-level state, and a spin-locking measurement is described by the sequence of three consecutive pulses

$$(-\pi/2)_Y - SL_X - (-\pi/2)_Y.$$  

The first pulse rotates the spin around the y-axis by a $-\pi/2$ angle resulting in the spin oriented along the x-axis. The second pulse — a so-called spin-locking pulse — is a long pulse with the variable amplitude and duration which is applied along the x-axis. Driven by the spin-locking pulse, the spin precesses around the x-axis at the Rabi frequency $\Omega_R$ (determined by the pulse amplitude), and, thus, the spin is effectively “locked” along the x-axis. The third pulse aligns the spin along the z-axis which allows one to measure the qubit state (e.g., by dispersive readout). In spin-locking measurements, both the amplitude and duration $\tau$ of the spin-locking pulse are varied, and, at a given amplitude, the relaxation rate is extracted from an exponential fit of the qubit excited-state population decay, $P_e = (1 + \exp(-\Gamma_1 \tau))/2$. According to the model of generalized Bloch equations (GBE) [3, 10, 11], the relaxation rate $\Gamma_1$ and relaxation time $T_{1\rho}$ of the qubit driven-state in the rotating frame (hence, the subscript symbol “$\rho$”) is used conventionally are given by:

$$\Gamma_1 = \frac{1}{T_{1\rho}} = \frac{1}{2} \Gamma'_1 + \Gamma_1(\Omega_R),$$

where the relaxation rate $\Gamma_1(\Omega_R)$ is related to the low-frequency longitudinal noise at the Rabi frequency $\Omega_R$, while the relaxation rate $\Gamma'_1$ is usually assumed to be determined by the high-frequency transverse noise at the qubit frequency $\omega_q$ ($\omega_q \gg \Omega_R$) [3, 10, 11]:

$$\Gamma'_1 \approx \Gamma_1(\omega_q) = (T_1(\omega_q))^{-1},$$

where $T_1$ is the qubit energy-relaxation time. However, it is known that the exact equation for $\Gamma'_1$ is given by [10, 11, 12, 13]:

$$\Gamma'_1 = \frac{1}{2} \left( \Gamma_1(\omega_q - \Omega_R) + \Gamma_1(\omega_q + \Omega_R) \right),$$

and, therefore, Eq. (2) is valid only if $\Gamma_1(\omega_q) \approx \Gamma_1(\omega_q \pm \Omega_R)$. Equation (3) is usually disregarded, and results of spin-locking measurements are interpreted as the characterization of noise sources at low frequencies determined by the Rabi frequency, since the term $\Gamma'_1/2 \approx \Gamma_1(\omega_q)/2$ is not expected to depend on $\Omega_R$, and, moreover, the condition $\Gamma_1(\Omega_R) \gg \Gamma_1(\omega_q)/2$ is fulfilled in many cases. For example, spectral features observed in the MHz frequency range in spin-locking experiments with a superconducting qubit were attributed to low-frequency fluctuators [10].

In this Rapid Communication, we present results of spin-locking measurements of a superconducting flux qubit coupled to off-resonant high-frequency two-level-system (TLS) defects. We observed spectral features in...
the driven-state relaxation time $T_{1p}$, which were caused by the interaction between the qubit with the frequency $\omega_q$ and defects with the frequencies $\omega_{\text{TLS}}$ matching the conditions $\omega_{\text{TLS}} = \omega_q - \Omega_R$ or $\omega_{\text{TLS}} = \omega_q + \Omega_R$. Thus, both low-frequency and high-frequency noise sources should be taken into account when interpreting results of spin-locking noise-spectroscopy measurements.

Our sample is a capacitively-shunted (c-shunt) superconducting flux qubit embedded in a 3D microwave cavity. Detailed information about the qubit design and experimental setup can be found in Ref. [16]. In the reported spin-locking experiments, a commercial high-power microwave amplifier from Spacek Labs with a gain of 40 dB was used to drive the qubit. The qubit was measured by dispersive readout. The cavity frequency was $\omega_c/2\pi \approx 8.2185$ GHz, the cavity linewidth was $\kappa/2\pi \approx 1$ MHz, and the dispersive frequency pull was $2\chi/2\pi \approx 1.4$ MHz. The qubit frequency measured as a function of the applied magnetic flux $\Phi$ is shown in Fig. 1(a).

FIG. 1. (a) The qubit frequency as a function of the applied magnetic flux. Dashed lines correspond to positions of TLS defects determined from $T_1$ measurements. (b) The energy-relaxation time $T_1$ of the qubit. Markers show approximate positions of four TLS defects resonantly coupled to the qubit. (c) The frequency of Rabi oscillations at a given drive amplitude. (d) The frequency of Rabi oscillations at the optimal bias point $\Phi = 0.5\Phi_0$ as a function of the drive amplitude. The black dashed line corresponds to a linear fit.

Variations of the $T_1$ values in the range 50–90 $\mu$s were related to the temporal variations reported previously [16] which were caused by quasiparticle tunneling [8, 17, 18] and a bath of background TLS defects [19–21]. In a separate two-tone measurement, the qubit anharmonicity was found to be about 0.8 GHz. At the optimal bias point, qubit decoherence times $T_{2E} \approx 6\mu$s and $T_{2R} \approx 5\mu$s were measured using spin-echo and Ramsey pulse sequences, respectively. In comparison with previous measurements of the sample [16], the qubit frequency was slightly lower due to the thermal cycling between the experimental runs, and the reduction of decoherence times $T_{2E}$ and $T_{2R}$ was caused by the low-frequency noise of the high-power amplifier.

The magnetic-flux dependence of the qubit energy-relaxation time $T_1$ is shown in Fig. 1(b). Variations of the $T_1$ values in the range 50–90 $\mu$s were related to the temporal variations reported previously [16] which were caused by quasiparticle tunneling [8, 17, 18] and a bath of background TLS defects [19–21]. In addition, we observed four pronounced dips which can be attributed to a resonant coupling between the qubit and a subset of distinct TLS defects denoted as TLS-1, TLS-2, TLS-3, and TLS-4. The dependence of the energy-relaxation rate on parameters of a coupled qubit-defect system is non-trivial [22–24]. Qualitatively, the presence of observed TLS defects did not cause averaged crossings in the qubit spectrum, and, hence, the qubit-defect couplings were weak which was in contrast with the previous works on strongly-coupled qubit-defect systems [3, 23, 29].

Figure 1(c) shows the frequency of Rabi oscillations at a given drive amplitude as a function of the applied magnetic flux. The decrease of the Rabi frequency at $\Phi \approx 0.5\Phi_0$ was probably due to the coupling to the bath of high-frequency TLS defects. The dependence of the Rabi frequency on the drive amplitude at the optimal bias point $\Phi = 0.5\Phi_0$ is shown in Fig. 1(d). The non-linear behavior of $1/T_1$ at high amplitudes was caused by the operation of the high-power microwave amplifier.

Results of spin-locking measurements at the optimal bias point $\Phi = 0.5\Phi_0$ are shown in Fig. 2. We used two types of spin-locking pulse sequences. The first one (labeled ‘S1’) was the standard sequence $(-\pi/2)Y - \text{SL}_X - (-\pi/2)Y$, in which the qubit was in the state $|0\rangle_x = (|0\rangle + |1\rangle)/\sqrt{2}$ (parallel to the $x$-axis) at the end of the first pulse. Here, the states $|0\rangle$ and $|1\rangle$ are eigenstates of the qubit Pauli operator $\sigma_x$ with eigenvalues +1 and −1 respectively. In experiments, we measured the probability to find the qubit in the excited state $|0\rangle$ in the end of the spin-locking sequence. Figure 2(a) shows the data set $D_1$ obtained using the S1 sequence. The second sequence (labeled ‘S2’) was the modified spin-locking sequence $(+\pi/2)Y - \text{SL}_X - (+\pi/2)Y$, in which the qubit was in the state $|1\rangle_x = (|0\rangle - |1\rangle)/\sqrt{2}$ (anti-parallel to the $x$-axis) at the end of the first pulse. The data set $D_1$ obtained using the S2 sequence is shown in Fig. 2(b). For both types of pulse sequences, rectangular pulses were used. The $\pi/2$-pulse duration was chosen to be long ($\approx 100$ ns) to minimize spectral widths of the $\pi/2$-pulses. The amplitude and duration of $\text{SL}_X$-pulses were varied. Figure 2(c) shows the arithmetic mean $P = (D_1 + D_2)/2$ which was used for estimations of the relaxation time $T_{1p}$.

A distinct spectral feature was observed in spin-locking measurements at the Rabi frequency $\Omega_R/2\pi \approx 51.3$ MHz [Figs. 2(a)–(c)]. We suppose that it was caused by the coupling between the qubit and the TLS-3 defect. Following Refs. [10, 24, 27, 30], we model the
FIG. 2. Results of spin-locking spectroscopy at the optimal point \( \Phi = 0.5\Phi_0 \). (a) The data set \( D_1 \) represents results obtained using the pulse sequence \((-\frac{1}{2})_Y - SL_X - (-\frac{1}{2})_Y\). (b) The data set \( D_2 \) corresponds to results obtained using the pulse sequence \((\frac{1}{2})_Y - SL_X - (+\frac{1}{2})_Y\). (c) The arithmetic \((D_1 + D_2)/2\) is the microwave drive term. Here, \( \sigma \) is the interaction term, and \((D_1, D_2)\) correspond to the frequency of the applied microwave drive.

The qubit-TLS system subjected to a microwave excitation by the Hamiltonian \( H = H_q + H_{TLS} + H_{1} + H_{MW} \), where \( H_q = (\hbar \omega_q/2)\sigma_z^{(1)} \) is the qubit Hamiltonian, \( H_{TLS} = (\hbar \omega_{TLS}/2)\sigma_z^{(2)} \) is the defect Hamiltonian, \( H_1 = \hbar g\sigma_x^{(1)} \sigma_z^{(2)} \) is the interaction term, and \( H_{MW} = \Omega_R \cos(\omega_{MW})\sigma_x^{(1)} \) is the microwave drive term. Here, \( \sigma_{x,y,z}^{(1,2)} \) are Pauli operators for the qubit (\( \sigma_x \), \( \sigma_y \), \( \sigma_z \)) and the coupling strength, and \( \omega_{MW} \) corresponds to the frequency of the applied microwave drive. In the case of the resonant driving (\( \omega_{q} = \omega_{MW} \)), the Hamiltonian can be written in the rotating wave approximation:

\[
H_{R} = \frac{\Omega_R}{2} \sigma_z^{(1)} + \frac{\Delta_{TLS}}{2} \sigma_z^{(2)} + g(\sigma_x^{(1)} \sigma_x^{(2)} + \sigma_y^{(1)} \sigma_y^{(2)})/2,
\]

where \( \Delta_{TLS} = \omega_{TLS} - \omega_{q}, \) and \( \sigma_{x,y,z}^{(1,2)} = (\sigma_{x,y,z}^{(1,2)} \pm i\sigma_{y,x,z}^{(1,2)})/2 \). Using QuTiP package \([31]\), we numerically simulated the dynamics of the system in the rotating frame by solving the Lindblad master equation for given energy-relaxation rates \( \Gamma_{1}^{(1)} \) and \( \Gamma_{1}^{(2)} \), and dephasing rates \( \Gamma_{2}^{(1)} \) and \( \Gamma_{2}^{(2)} \) of the qubit and the TLS defect, respectively (with the corresponding collapse operators \( \sqrt{\Gamma_{1}^{(1)}} \sigma_{+}^{(1)}, \sqrt{\Gamma_{1}^{(2)}} \sigma_{-}^{(2)}, \sqrt{\Gamma_{2}^{(1)}}/2 \sigma_{+}^{(1)}, \) and \( \sqrt{\Gamma_{2}^{(2)}}/2 \sigma_{-}^{(2)} \)). Following Ref. \([32]\), we used the values \( \Gamma_{1}^{(2)} = 10^6 s^{-1} \) and \( \Gamma_{2}^{(2)} = 0 \) for the defect relaxation rates, and, for simplicity, we assumed the qubit was free of pure dephasing, \( \Gamma_{2}^{(1)} = 0 \). To simulate the S1 (S2) measurement, the defect was initialized in its ground state in the laboratory frame, and the expectation value \( \langle \sigma_x^{(1)} \rangle \) as a function of the evolution time was calculated for the qubit initial state \( |0\rangle_x \) (\( |1\rangle_x \)) and converted to the corresponding value of \( D_1 = (1 + \langle \sigma_x^{(1)} \rangle)/2 \) (\( D_1 = (1 - \langle \sigma_x^{(2)} \rangle)/2 \)). We found that the observed spectral feature can be well reproduced using the following parameters: \( \Omega_R/2\pi = \Delta_{TLS}/2\pi = 51.3 \text{MHz} \), the coupling strength \( g/2\pi = 28 \text{kHz} \), and the qubit energy-relaxation rate \( \Gamma_{1}^{(1)} = 1.5 \times 10^6 s^{-1} \) [Fig.2(d)]. Thus, conditions of the Purcell-like effect were fulfilled in the rotating frame: \( \Gamma_{1}^{(1)} < g < \Gamma_{1}^{(2)} \). We assume that the qubit effective dissipation rate can be roughly estimated as \( \Gamma_P \approx g^2/\Gamma_{1}^{(2)} \approx 0.31 \times 10^5 s^{-1} \) extracted from an exponential fit of the spin-locking data P [Fig.2(d)]. Therefore, the phase cycling procedure described above — averaging the data sets obtained using S1 and S2 sequences with alternating pulse phases — is a valid method to obtain the information on the driven-state relaxation rate. The \( T_{1p} \) data, presented in Fig.3(a), was extracted from the signal P using exponential fitting.

Besides the main feature at \( \Omega_R/2\pi = 51.3 \text{MHz} \), various additional spectral features were observed at other Rabi frequencies [Fig.2(a)–(c) and Fig.3(b)]. All features can be divided into two groups based on their “polarity” that is defined by the relation between the steady levels of the excited-state population \( D_1 \) and \( D_2 \) in S1 and S2 measurements, respectively. For “positive polarity” features, such as the one observed at \( \Omega_R/2\pi = 70 \text{MHz} \) in Figs.2(a),(b), the condition \( D_1 > D_2 \) was fulfilled. For those features, steady values of \( \langle \sigma_x^{(1)} \rangle \) in the end of the spin-locking pulse were positive \([35]\), which was caused by the heating by defects with \( \Delta_{TLS} < 0 \) according to our model \([30]\). Other features, such as the main one observed at \( \Omega_R/2\pi = 51.3 \text{MHz} \), can be attributed to the “negative polarity” group for which the condition \( D_1 < D_2 \) was realized. For those features, steady values of \( \langle \sigma_x^{(1)} \rangle \) were negative, which was caused by the cooling of the qubit by defects with \( \Delta_{TLS} > 0 \). It should be mentioned that the pure qubit state \( |0\rangle_x \) (\( |1\rangle_x \)) corresponded to the excited (ground) state in the rotating frame with \( \langle \sigma_x^{(1)} \rangle = 1 \) (\( \langle \sigma_x^{(1)} \rangle = -1 \)), and, therefore, the steady spin-locking state of the qubit without coupling to the TLS defect would be a completely mixed state of \( |0\rangle_x \) and \( |1\rangle_x \) states with \( \langle \sigma_x^{(1)} \rangle = 0 \).

Figure3(a) shows relaxation times \( T_{1p} \) as a function of the Rabi frequency at the optimal bias point and the flux bias detuned from the optimal point by \( \delta \Phi = \Phi - 0.5\Phi_0 \approx 7.5 \times 10^{-4} \Phi_0 \). At the detuned bias, the qubit frequency was 4.33 GHz, and, hence, the qubit frequency detuning was \( \delta\omega_q/2\pi \approx 3 \text{MHz} \). In addition, Fig.3(a) shows the values of \( (\Gamma_{1}^{(2)}/2)^{-1} \) at the op-
The relaxation rate $\Gamma'_1(\Omega_R)$ was estimated from the data shown in Figs. 2(a), (b) using the equation $\Gamma'_1 \approx \Gamma_1(\omega_q) / 2 + \Gamma_1(\omega_q + \Omega_R) / 2$ obtained from Eq. 4 under the assumption $\Gamma_1(\omega_q - \Omega_R) \approx \Gamma_1(\omega_q)$ (here, we neglected the contribution from defects with $\Delta_{TLS} < 0$). At low Rabi frequencies, in accordance with Eq. 1, the $T_{1p}$ time was dominated by the term $\Gamma_1^{-1}(\Omega_R)$ due to the low-frequency $1/f$ noise of the high-power amplifier. At high Rabi frequencies, the relaxation time plateaued at the value of 100 $\mu$s which was below the baseline level of $(\Gamma'_1/2)^{-1}$. In this frequency range, the $1/f$ noise was not dominant, and the difference between background values of $T_{1p}$ and $(\Gamma'_1/2)^{-1}$ was caused by other low-frequency noises at $\Omega_R$, and by high-frequency TLS defects with $\Delta_{TLS} < 0$ and short relaxation times (broad spectra). The pronounced features in the $T_{1p}$ data were caused by the coupling between the qubit and high-frequency TLS defects with long relaxation times (narrow lines). As explained above, the main feature observed at the optimal point was attributed to the TLS-3 defect. At the detuned flux bias, positions of spectral features related to the defects with $\Delta_{TLS} > 0$ were shifted to lower Rabi frequencies by the value of the qubit frequency detuning $\delta \omega_q$. Therefore, we identified the feature observed at $\Omega_R/2\pi \approx 45$ MHz at the detuned bias as the main one corresponding to the TLS-3 defect [Fig. 3(a)]. As shown in Fig. 3(b), there were some discrepancies in the positions of TLS defects observed in $T_1$ and $T_{1p}$ measurements which could be caused by fluctuations of defect frequencies [10, 21]. In separate spin-locking measurements, we found that positions of the spectral features were not affected by in-plane magnetic fields of up to 0.2 mT, which indicated that the detected TLS defects were charge defects.

In conclusion, we demonstrated that spin-locking measurements of the driven-state relaxation of a superconducting qubit with the frequency $\omega_q$ can be significantly affected by the interaction with off-resonant TLS defects with the frequencies $\omega_{TLS}$ if one of the conditions $\omega_{TLS} = \omega_q \pm \Omega_R$ is met. Although the qubit and defects were nominally off-resonance in the laboratory frame in our experiments, the Purcell-like regime of resonant qubit-defect coupling was realized in the rotating frame by driving the qubit with the spin-locking pulse. As a result, the qubit relaxation was affected by the interaction with the defects. Thus, in addition to the effect of low-frequency noise sources reported previously [10], spectral features in spin-locking measurements can be caused by off-resonant high-frequency defects. Similar effects can be observed in other types of driven-state noise spectroscopy such as Rabi and rotary-echo measurements [4, 7, 8]. The reported qubit-defect coupling can be also interpreted in terms of Autler-Townes splitting of the qubit state as the interaction between the defect and one of the qubit dressed states [37, 38]. Our results demonstrate that spin-locking methods can be used to couple a superconducting qubit to other quantum systems with high-frequency transitions such as NV centers in diamond [29]. Our work also shows that spin-locking techniques can be used to determine the spectral distribution of high-frequency defects in the vicinity of the qubit frequency, which can be especially useful for fixed-frequency qubits such as transmons.

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At a given Rabi frequency, the steady levels of the expectation value \( \langle J_z \rangle \) in the end of the spin-locking pulse were the same for S1 and S2 measurements.

For defects with \( \Delta_{\text{TLS}} < 0 \), the ground state in the laboratory frame corresponds to the excited state in the rotating frame. Therefore, such defects can be considered as “hot” defects.
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