Implementation of the three-qubit phase-flip error correction code with superconducting qubits

Downloaded from: https://research.chalmers.se, 2020-04-25 06:21 UTC

Citation for the original published paper (version of record):
Implementation of the three-qubit phase-flip error correction code with superconducting qubits
Physical Review B - Condensed Matter and Materials Physics, 77
http://dx.doi.org/10.1103/PhysRevB.77.214528

N.B. When citing this work, cite the original published paper.
Implementation of the three-qubit phase-flip error correction code with superconducting qubits

L. Tornberg,1 M. Wallquist,2 G. Johansson,1 V. S. Shumeiko,1 and G. Wendin1
1Chalmers University of Technology, SE-41296 Gothenburg, Sweden
2Institute for Theoretical Physics, University of Innsbruck, and Institute for Quantum Optics and Quantum Information
of the Austrian Academy of Sciences, 6020 Innsbruck, Austria
(Received 11 March 2008; published 27 June 2008)

We investigate the performance of a three-qubit error correcting code in the framework of superconducting qubit implementations. Such a code can recover a quantum state perfectly in the case of dephasing errors but only in situations where the dephasing rate is low. Numerical studies in previous work have however shown that the code does increase the fidelity of the encoded state even in the presence of high error probability, during both storage and processing. In this work we give analytical expressions for the fidelity of such a code.

In the regime of fast gate operations, the effects of correlated noise have also been studied.20 In this paper we take a different approach, where all single-qubit operations are considered to be much faster than the dephasing time \( t_{\text{deph}} \ll T_2 \). We can therefore neglect the errors that occur on this time scale. As the two-qubit gate we consider the controlled-phase (cPhase) gate which is diagonal in the computational basis. Two realizations of this gate, relevant for superconducting implementations, are considered. In the first case, the coupling Hamiltonian itself is diagonal (experimentally achievable via e.g., a large Josephson junction21). This allows us to calculate the fidelity of the corrected state analytically. In the second case the qubit-qubit coupling is mediated by a cavity bus22 where the dipole coupling between the cavity and quantum systems have recently attracted much attention due to their long coherence times.23,24 In this case the coupling is however not diagonal and we have to calculate the fidelity of the QECC numerically. We show that the fidelity of the QECC when using the cavity-mediated coupling is comparable to the case when the diagonal cPhase gate is used. The structure of the paper is as follows. In Sec. II we briefly discuss the model used to describe dephasing and introduce the three-qubit QECC with instantaneous gates. In Secs. III and IV we consider realistic implementations of gates. The diagonal coupling is considered in Sec. III and the cavity-mediated coupling in Sec. IV. We conclude in Sec. V.

II. INSTANTANEOUS GATES

We begin by describing the QECC in the ideal case, where only single errors are present, and the gates are instantaneous. In Sec. II B we relax the first approximation and study the case of multiple errors.

The Hamiltonian for a qubit coupled longitudinally to a heat bath of harmonic oscillators is given by

\[
H = -\frac{E}{2} \sigma_z + \sigma_i \sum_{l} \hbar \lambda_{il} (b_{il}^\dagger + b_{il}) + \sum_{i} \hbar \omega_i b_{il}^\dagger b_{il},
\]  

(1)

where \( E \) is the qubit level splitting and \( b_{il}^\dagger \) and \( b_{il} \) the usual creation/annihilation operators of the harmonic oscillator.
With this Hamiltonian, one can derive a master equation in the Markov limit\cite{25}
\[ \dot{\rho}(t) = -\frac{\Gamma_e}{2}(\rho(t) - \sigma_e \rho(t) \sigma_e), \]  
where $\Gamma_e$ is the dephasing rate given by $\Gamma_e = 4 \text{ Re } L(0)$, where $\text{Re } L(\omega)$ is the real part of the correlation function (noise spectral density) of the bath
\[ \text{Re } L(\omega) = \pi \eta(\omega) \lambda^2(\omega) \delta(\omega, T) \times \left[ \exp\left( \frac{i \hbar \lambda(\omega)}{k_B T} \right) \Theta(\omega) + \Theta(-\omega) \right]. \]

Here $\hbar \lambda(\omega)$ is the coupling energy to the bath, $\eta(\omega)$ is the bath density of states, $\delta = [\exp(i \hbar \omega / k_B T) - 1]^{-1}$ the Bose occupation number, and $\Theta(x)$ is the step function where we adopt the convention $\Theta(0) = 1/2$. We leave these quantities unspecified since the experimental parameter of interest is the dephasing rate itself. (See Appendix A for further details). Solving Eq. (2) gives an exponential decay of the off-diagonal elements in the density matrix
\[ \rho_{ij} = e^{-\Gamma_e t} \rho_{ij}(0). \]

The diagonal elements are stationary in time. The results for the diagonal and off-diagonal elements can be combined into the solution
\[ p(t) = p(0) + [1 - \exp(-\Gamma_e t)] \sigma_e, \]
where $p(t)$ can be given the meaning of a time-dependent probability for the qubit to be in the correct state
\[ p(t) = \frac{1}{2} [1 + \exp(-\Gamma_e t)]. \]

In this paper we assume the dephasing to act independently on each qubit. In the presence of noise all three qubits in the circuit will thus evolve according to Eq. (5). This form allows us to interpret dephasing as a stochastic process where a phase flip occurs with probability $1 - p(t)$. This discretization of the continuous phase damping is the very core of quantum error correction.

A. One qubit error

The three-qubit QECC that corrects for phase-flip errors is depicted in Fig. 1. The qubit is initially in the state $|\psi\rangle = \alpha |g\rangle + \beta |e\rangle$ and the ancillas are initiated in the ground state $|gg\rangle$. The first part of the circuit (dashed box) serves to entangle the three qubits leaving them in the encoded state
\[ |\Psi\rangle = \alpha |ggg\rangle + \beta |eee\rangle. \]

The role of the Hadamard gates is to rotate the phase flip of Eq. (5) into a bit flip $H \sigma_e H = \sigma_e$. Thus, during the part of the evolution where errors occur, all three qubits will evolve according to Eq. (5), but with $\sigma_e$ replaced with $\sigma_z$. If the dephasing is weak $\Gamma_e t \ll 1$, the error probability is small $1 - p(t) \ll 1$, and we can neglect terms representing errors occurring on more than one qubit at a time. Just before the decoding part (dotted box), the three qubits are in a mixed state with the following probabilities to find the system in the respective pure states:
\[ p^3: |a|ggg\rangle + |b|eee\rangle, \]
\[ p^2(1 - p): |a|ggg\rangle + |b|eee\rangle, \]
\[ p^2(1 - p): |a|ggg\rangle + |b|eee\rangle, \]
\[ p^2 (1 - p): |a|ggg\rangle + |b|eee\rangle. \]

After the decoding the states are given by
\[ p^3: (|a|g\rangle + |b|e\rangle)|gg\rangle, \]
\[ p^2(1 - p): (|a|e\rangle + |b|g\rangle)|ee\rangle, \]
\[ p^2(1 - p): (|a|g\rangle + |b|e\rangle)|eg\rangle, \]
\[ p^2 (1 - p): (|a|g\rangle + |b|e\rangle)|ge\rangle. \]

This discretization of the continuous phase damping is the very core of quantum error correction.

FIG. 1. The three-qubit QECC for correcting phase-flip errors. The state of the information-carrying qubit is encoded with three qubits, using two ancilla qubits which are initially in the state $|g\rangle$. After decoding, a single phase flip on the information-carrying qubit is corrected by a Toffoli gate, controlled by the two ancilla qubits. A single-qubit operation, conditioned by a measurement of the two ancilla qubits, can substitute for the Toffoli gate.
Thus in principle we can correct repeatedly and derive an effective dephasing rate. The probability $p_{c,n}$ to be in the correct state after $n$ such correction cycles obeys the equation

$$
\begin{align*}
\left( p_{c,n} \right) = \left( \begin{array}{cc}
1 - p_{c,t} & p_{c,t} \\
1 - p_{c,t} & p_{c,t}
\end{array} \right)^n \left( \begin{array}{c}
1 \\
0
\end{array} \right),
\end{align*}
$$

(12)

where $(0,1)$ and $(0,1)$ denote the correct and flipped state, respectively, and $t_c$ is the duration of the cycle. This equation is easily solved yielding

$$
p_{c,n} = \frac{1}{2} \left[ 1 + \left( \frac{3 - e^{-2 \Gamma \varphi \ell_c}}{2} \right) e^{-\Gamma \varphi \ell_c} \right].
$$

(13)

Comparing Eqs. (6) and (13) with $t=t_n$, we may derive an effective dephasing rate $\Gamma_{\text{eff}}$ such that $p_{c,n}(t) = \frac{1}{2} \left[ 1 + e^{-\Gamma_{\text{eff}} \ell_c} \right]$. The effective rate $\Gamma_{\text{eff}}$ is related to $t_c$ and $\Gamma_c$ according to

$$
\Gamma_{\text{eff}} = \Gamma_c \left[ 1 - \frac{1}{2 \Gamma_c} \left( \frac{3 - e^{-2 \Gamma \varphi \ell_c}}{2} \right) \right],
$$

(14)

giving $\Gamma_{\text{eff}} = \Gamma_c$ for all $\Gamma_c$, as can be seen in Fig. 2 where the ratio $\Gamma_{\text{eff}}/\Gamma_c$ is plotted as a function of $\Gamma_c$. We also indicate the value of $\Gamma_{\text{eff}}/\Gamma_c$ that can be achieved with current examples of superconducting qubit implementations. Here we have estimated the repetition time as $t_c \sim h/g$, $g$ being the qubit-qubit coupling energy. We see that, even for realistic parameters, there can be a significant increase in coherence in the encoded quantum state. This shows that the three-qubit QECC can be used to prolong the lifetime of the qubit even when the assumption of instantaneous gates is relaxed.

### III. DIRECT QUBIT COUPLING

We now depart from the approximation of perfect processing and consider the regime where $\Gamma_{\text{aq}} \ll 1$, taking into account the errors that occur during the two-qubit gate operations. We however make the assumption of fast single-qubit operations $t_{\text{aq}} \ll 1/\Gamma_c$ and neglect any dephasing occurring on this time scale. We first consider the performance of the QECC protocol when the information-carrying qubit is coupled directly to each ancilla qubit with a coupling which is diagonal in the energy eigenbasis

$$
H = \sum_i E_i \sigma_i^{(i)} + \lambda_i \sigma_z^{(i)} \sigma_z^{(i+1)},
$$

(15)

Such a coupling can be realized using e.g., the circulating currents in ring-shaped CPB qubits$^{30}$ interacting via a large Josephson junction.$^{21}$ This coupling, together with single-qubit phase gates naturally gives rise to the general controlled-phase gate: $\text{diag}[1,1,1,\exp(i 4 \pi \ell_c \hbar)]$, which for the interaction time $\lambda T = h \pi/4$ generates the cPhase gate. The advantages of studying the implementation of QECC with this setup first are: (1) the cPhase gate is also a natural gate for the cavity-mediated qubit coupling, and (2) using the cPhase gate to implement the cNOT gates as shown in Fig. 3, the error operators $I$ and $\sigma_z$ [see Eq. (5)] commute with all gates inside the circuit except for the single-qubit Hadamard gates. This implies that the errors occurring during the ex-
execution of the cPhase can be moved to the end or beginning of the gate as indicated in Fig. 4. Denoting the gate operation time \( t_g \) and the storage time \( t_s \), we see that the ancillas are subject to the same errors as before, but now during a longer time \( 2t_g + t_s \). The error on the uppermost line, however, must be divided into three parts, since it is separated by the Hadamard gates. Calculating \( r(t) \) thus reduces to the calculation of \( 2^{-8} \cdot 2 = 32 \) terms in the operator sum representation, which is analytically tractable (cf. 8.8.8.8 = 512 terms if the errors on the ancillas could not be collected). Since the Toffoli gate can be replaced with fast measurements we also neglect errors occurring during the execution of this. The final state of the qubit after detection and correction is given by the density matrix

\[
\rho(t) = p_1 \rho_1 + p_2 \rho_2 + p_3 \rho_3 + p_4 \rho_4,
\]

with

\[
\rho_1 = 1, \rho_2 = P_x \rho(t_g)^2 + [1 - p(t_g)],
\]

\[
\rho_3 = \sigma_y \rho \sigma_y, \rho_4 = 2 P_x \rho(t_g)[1 - p(t_g)],
\]

\[
\rho_5 = \sigma_z \rho \sigma_z, \rho_6 = 2 P_x \rho(t_g)[1 - p(t_g)],
\]

\[
\rho_7 = \sigma_z \rho \sigma_z, \rho_8 = 2 P_x \rho(t_g)[1 - p(t_g)],
\]

with \( p(t_g) \) as in Eq. (6) and the probabilities \( P_0 \) and \( P_1 \) given by

\[
P_0 = \frac{1}{4} \left( 1 - e^{-i \theta} \right) + \frac{1}{2} e^{-i \theta} + \frac{1}{4} e^{i \theta}.
\]

Since there now is more than one type of error present in the final state of the qubit, it is no longer meaningful to compare the error probabilities directly. Instead we use the fidelity

\[
F(t) = \langle \psi | \rho(t) | \psi \rangle \tag{18}
\]

between the initial state \( | \psi \rangle \) and the corrected state \( \rho(t) \) of Eq. (16) to quantify the benefit of using the QECC. Since \( F(t) \) depends on the initial state \( | \psi \rangle \) we use the minimum fidelity as a measure of the code performance

\[
F_{\text{min}} = \min_{| \psi \rangle} F(t). \tag{19}
\]

To find this minimum fidelity, we make the observation that \( P_s \geq P_0 \) for all times \( t_g, t_s \). This, together with the fact that \( p(t_g)^2 + [1 - p(t_g)]^2 \geq 2 p(t_g)[1 - p(t_g)] \) for all \( t_g \) gives us the inequalities [see Eq. (16)] \( p_1 \geq p_2 \geq p_3 \) and \( p_4 \geq p_5 \). This in turn gives the minimum fidelity according to

\[
F_{\text{min}} = \min \left( F(t) \right).
\]

We now compare the minimum fidelity between the state in Eq. (16) and the uncorrected state which is subject to the same dephasing given by

\[
F_{\text{uncorr}} = p(t),
\]

with \( t = 2t_g + t_s \). The difference in fidelities \( F_{\text{min}} - F_{\text{uncorr}} \) is plotted in Fig. 5. For perfect gates, coding improved the fidelity for all times. The situation is drastically different when errors occur during processing. There is now a lower limit on the speed of the gate operations, given by \( F_{\text{min}} = F_{\text{diag}} \) above which QECC is beneficial. This condition can equally be stated as a relation between \( t_g \) and \( t_s \) given by

\[
t_g = \frac{1}{2 \Gamma} \left( -6 \Gamma \psi t_g - \log \left( e^{-\frac{2 \Gamma \psi t_g}{3}} \right) \right).
\]

For typical values of the dephasing time in superconducting systems \( t_g \approx 1 \mu s \), this gives a maximum gate operation time of the order 0.1 \( \mu s \). From Fig. 5 it is clear that, for each fixed value of \( \Gamma \psi t_g \), there exists an optimum storage time such that the difference \( F_{\text{min}} - F_{\text{diag}} \) is maximized

\[
t_f^{\text{opt}} = -\frac{1}{2 \Gamma \psi} \left[ 2 \Gamma \psi t_f + \log \left( \frac{2 - e^{2 \Gamma \psi t_f}}{3} \right) \right].
\]

This is plotted with a solid black line in the \( t_f t_g \) plane of Fig. 5.
We can understand the origin of the optimum in the following way: For instantaneous gates $t_s=0$, we recover the result of Sec. II B where coding is beneficial for all finite storage times $t_s$. Along the line of zero storage time $t_s=0$, however, we see there is nothing to be gained from using the code. This is simply because the logical qubit is only shielded when it is encoded, i.e., during the storage. For zero storage time, we are only applying a series of faulty gates, which of course reduces the fidelity. As $t_s$ grows, there is a competition between these uncorrectable errors and the beneficial effects of the code obtained for any finite $t_s$. As $t_s$ exceeds the limit given in Eq. (23) the positive effects of the code can however no longer compensate for the uncorrectable errors in the gates. Finally, as $t_s \to \infty$, there is no longer any advantage in using the code and the fidelity approaches that of the unprotected qubit.

We now consider what happens with the condition given in Eq. (23) when the code is repeated $n$ times, each run taking time $t_{s}^{ge}=2t_s+t_g$. The four probabilities $p_{i}^{n}$ describing the state of the qubit after $n$ such cycles obey the equation

$$
\begin{pmatrix}
    p_{1,n} \\
    p_{2,n} \\
    p_{3,n} \\
    p_{4,n}
\end{pmatrix} =
\begin{pmatrix}
    p_1 & p_z & p_x & p_y \\
    p_z & p_1 & p_y & p_x \\
    p_x & p_y & p_1 & p_z \\
    p_y & p_x & p_z & p_1
\end{pmatrix}^n
\begin{pmatrix}
    1 \\
    0 \\
    0 \\
    0
\end{pmatrix},
$$

in the basis $\{p_1, p_2, p_3, p_4\}$. It is useful to express the probabilities $p_i^n$ in term of the eigenvalues of the $P_{ge}$ matrix

$$p_{1,n} = \frac{1}{4} (1 + \lambda_2^n)(1 + \lambda_3^n),$$

where the eigenvalues are given by

$$\lambda_1 = 1,$$

$$\lambda_2 = 1 - 2(p_x + p_y),$$

$$\lambda_3 = 1 - 2(p_x + p_y),$$

$$\lambda_4 = \lambda_2 \lambda_3.$$  

From Eq. (26) and the fact that $\lambda_i \geq 0$ for all times, it is clear that $p_{1,n} = p_{2,n}$. The minimum fidelity in Eq. (20) thus generalizes to the iterative case by replacing $p_1$ with $p_{1,n}$ and $p_s$ with $p_{2,n}$. Further the condition for the QEC to be beneficial, $F_{\text{diag}}^{\text{min}} \geq F_{\text{min}}^{\text{min}}$, in the iterative case gets the form

$$p_{1,n} + p_{2,n} = p(n t_{s}^{ge}),$$

which equivalently can be expressed as

$$\lambda_3 \lambda_4 \geq \tau_n,$$

where $\tau_n = 2p(n t_{s}^{ge}) - 1$ is the eigenvalue $\neq 1$ of the corresponding $P$ matrix for the case of no error correction. Since Eq. (29) is $n$ independent it holds for any $n$. In particular we can choose $n=1$ for which it coincides with the condition in Eq. (23). Since the code introduces additional errors this result is not obvious. We had rather expected a case where a combination of the different probabilities in Eq. (25) would result in a faster fidelity decay. The main conclusion from this section is that the three-qubit QEC can improve the fidelity of the qubit state, even in the case of errors during gate operations. In order to benefit from the code, one must however assure that the normalized gate operation time $\Gamma_{s}^{\text{eff}}$ can be made shorter than $\log 2/2 \approx 0.35$. We emphasize that the results in this section hold for a diagonal qubit coupling only.

**IV. COUPLING VIA TUNABLE CAVITY**

The system we consider for implementing the QEC is a set of Cooper-pair box (CPB) qubits which are capacitively connected to a one-dimensional superconducting stripline cavity. The strong-coupling regime has been achieved experimentally for this type of mesoscopic cavity-QED system. In this setup, the stripline is terminated with a superconducting quantum interference device (SQUID) whose effective inductance can be tuned by applying an external magnetic flux $\Phi_S$. Thus by changing $\Phi_S$ one changes the cavity boundary conditions which leads to a tunable resonance frequency $\omega_0(\Phi_S)$ of the cavity. The Hamiltonian of the decoupled cavity mode and the CPB qubits read, at the
oscillator, interrupted by idle periods

interference loops in the interaction between qubit 2 and the
cavity are decoupled as shown in Fig. 6.

excite higher oscillator states

\[ \frac{d}{dt} c_{\ell} = \frac{\hbar}{2} \sum_{i=1}^{3} E_{\ell i} a_{i}^\dagger a + \hbar \omega_c (\Phi_\ell) \left( a_{i}^\dagger a + \frac{1}{2} \right) - \hbar \omega_c (\Phi_\ell) \sum_{i=1}^{3} E_{\ell i} \sigma_i^z, \]

where \( E_{\ell i} \) is the Josephson energy of the \( i \)th qubit.

In Ref. 22 it has been shown how to perform the universal
two-qubit cPhase gate using the cavity as a bus mediating
the qubit interaction. The strategy is to tune the oscillator
through resonance with one qubit at a time, performing
\( \pi \)-pulse swaps in every step, and then reverse the sequence.
However, due to the photon number dependence of the
\( \pi \)-pulse swap operation time, \( T = \pi / g \sqrt{n + 1} \), one has to insert
interference loops in the interaction between qubit 2 and the
oscillator, interrupted by idle periods \( T_n \leq 1 / g \) where qubit
and cavity are decoupled as shown in Fig. 6.

Note that the oscillator must be tuned adiabatically not to
excite higher oscillator states \( \omega_\ell / \omega_c \ll 1 \).

In the implementation of the QECC protocol in Fig. 1,
two consecutive cNOT operations are used. The gate sequence
for the encoding is equivalent to two consecutive
cPhase gates in addition to single-qubit Hadamard gates, as
shown in Fig. 3. Let us focus on the operation of the cavity
and forget single-qubit gates for the moment since these are
performed independently of the cavity. In the first step of the
cPhase operation, the state of qubit 1 is swapped onto the
oscillator, \( (|g\rangle + |\beta e\rangle)|0\rangle \rightarrow (|g\rangle a_{i}^\dagger |0\rangle + |\beta e\rangle |1\rangle) \).
During the next steps the oscillator interacts with qubit 2 in order to create a
phase shift which depends on both the state of the oscillator
and the state of qubit 2. The last step is to swap the oscillator state back onto qubit 1: \( |g\rangle (a_{i}^\dagger |0\rangle + |\beta e\rangle |1\rangle) \rightarrow (|g\rangle a_{i}^\dagger |0\rangle - |\beta e\rangle |1\rangle) \)
(minus sign appears because of the form of the interaction
term). When we perform two cPhase gates after each
other, first involving qubits 1 and 2, then involving qubits 1
and 3, it is possible to shorten the sequence. Namely, at the
end of the first cPhase gate the state of the oscillator is
transferred to qubit 1, but then it is immediately transferred
back to the oscillator at the beginning of the second cPhase
sequence. At the very end, these two operations, \( |g\rangle (a_{i}^\dagger |0\rangle + |\beta e\rangle |1\rangle) \rightarrow (|g\rangle a_{i}^\dagger |0\rangle - |\beta e\rangle |1\rangle) \), correspond only
to a single-qubit rotation on qubit 1. Thus we can shorten
the encoding protocol by directly tuning the oscillator from qubit
2 to qubit 3, as shown in Fig. 7, and correcting the sign with
an additional diagonal single-qubit gate.

A. Simulating the three qubit+cavity system with dephasing

During the pulse sequence described in the previous
section, the qubits will interact with each other by entangling
their states with the oscillator. If dephasing is added to this
picture it is no longer tractable to solve the equation of motion
analytically to obtain a solution as that in Eq. (5). Instead
we simulate the system dynamics by solving its governing
master equation numerically. This section is divided
into two parts. In the first part we describe how to derive
master equation from a system plus bath interaction. In the
second part we discuss how this equation is implemented numerically.

The major difficulty with the analysis is that the
system Hamiltonian will depend on whether the qubit is on
or off resonance with the cavity. Because of this, we must be
careful when treating the coupling to the bath, which will
have different structure in the energy eigenbasis of the two
different cases. Another difference lies in the fact that we
have introduced an additional quantum system trough the
cavity. This is also subject to dephasing which must be taken
into account. The standard procedure is to couple the cavity
to a heat bath longitudinally through the cavity number
operator. Since only one qubit at a time is on resonance
with the cavity, our analysis can, without loss of generality,
be limited to the case of a single qubit plus cavity. The
Hamiltonian for the system plus baths will then be given by
\[ H = H_{\text{sys}} + \sigma X + nY + H_{\text{bath}}^X + H_{\text{bath}}^Y, \]  
where the baths are chosen to be a collection of harmonic oscillators and \( n = a^d a \) is the cavity number operator. In this case, \( H = \hbar \sum \delta (b_i + b_i^\dagger) \) and \( H_{\text{bath}}^X = \hbar \sum \delta \omega_n b_i b_i^\dagger \) with the oscillator bath defined analogously. The system Hamiltonian \( H_{\text{sys}} \) is given by

\[
H_{\text{sys}} = \begin{cases} 
H_0 = \hbar \omega_n \left( a^d a + \frac{1}{2} \right) - \frac{\hbar g}{2} \sigma_z, & \text{off resonance} \\
H_{JC} = H_0 + \frac{\hbar g}{2} (a^d \sigma_- - a \sigma_+), & \text{on resonance}
\end{cases}
\]

for the idle and resonant periods in the pulse sequence respectively.

1. Dephasing with qubit and cavity on resonance

In this section we show how to derive a master equation for the reduced density matrix of the qubit plus cavity system. For simplicity, we only consider the qubit-bath coupling. The cavity can be treated analogously and we refer to Appendix A for a detailed analysis. In the idle periods the system Hamiltonian commutes with the qubit-bath interaction. The situation is thus identical to that described by Eq. (2) which means that the coupling to the bath results in pure dephasing in the energy eigenbasis \( |g/e, n \rangle \)

\[
\rho_{g,e;n}(t) = \delta_{n0} e^{-\frac{\hbar g n}{2}} \rho_{g,e;n}(0).
\]

When the qubit is moved into resonance with the cavity the system is described by the JC Hamiltonian whose eigenstates are the dressed states with corresponding eigenenergies

\[
|n; \pm \rangle = \frac{1}{\sqrt{2}} \left( |g; n \rangle \pm i |e; n - 1 \rangle \right), \quad E_{n;\pm} = \hbar \omega_n \pm \frac{\hbar g n}{2},
\]

\[
|g; 0 \rangle, \quad E_{g0} = 0.
\]

What is important from the point of view of dephasing is that in the dressed state picture, the longitudinal coupling to the bath becomes transversal with matrix elements

\[
\langle 0; g | \sigma_3 | 0; g \rangle = 1, \quad \langle m; \pm | \sigma_3 | n; \pm \rangle = 0,
\]

\[
\langle 0; g | \sigma_3 | n; \pm \rangle = 0, \quad \langle m; \pm | \sigma_3 | n; \mp \rangle = \delta_{mn}.
\]

Hence in this basis the coupling will not induce pure dephasing, but will instead give rise to relaxation between the dressed states. We note that this operator is block diagonal which means that the relaxation will take place within the blocks of equal \( n \). Within such blocks the system behaves approximately like a two-level system coupled transversally to a bath with interaction Hamiltonian \( H = \sigma X \). Hence, the diagonal elements of the density matrix will approach equilibrium exponentially with the relaxation rate \( 1/T_1 \) given by the noise spectral density \( L(\omega) \) [see Eq. (3)], at the relevant transition frequencies given by \( \omega = g \sqrt{n} / \hbar \) as can be seen from Eqs. (35) and (36).

Similarly, coherences formed by superpositions of states within blocks of equal \( n \) will decay exponentially with a rate \( 1/T_2 = 1/2T_1 \). One important difference between the quasi-two-level systems formed by the dressed state blocks and a real two-level system is that the relaxation rates for different blocks will be different due to the nonharmonicity of the dressed state energy levels. In addition to this, coherences formed by superpositions of states between different blocks will have another set of decay rates. [This can be seen from the full Liouville equation for the reduced density matrix, see Eqs. (A1) and (A2).] However, we note that all relevant transition frequencies lie in a range given by the cavity-qubit coupling \( g \). Thus if we assume that the baths have no structure on this scale and that the temperature is much higher than the cavity-qubit coupling \( k_B T \gg g \), we can safely approximate the noise spectral density in Eq. (3) by its zero-frequency limit for all relevant transitions. The situation is similar for the oscillator-bath coupling, for which we make the same assumptions about the bath. Apart from setting all the rates equal this approximation has another important implication. Since all rates are taken at zero frequency, they will be the same for the resonant and off-resonant passages. With this clearly stated we can derive a master equation in the Markov approximation for the reduced cavity-qubit density matrix

\[
\dot{\rho} = -\frac{i}{\hbar} [H, \rho] - \gamma(n^2 \rho + \rho n^2 - 2n \rho n) - \frac{\kappa}{2} (\rho - \sigma \rho \sigma),
\]

with the rates \( \kappa \) and \( \gamma \) given in Appendix A along with a detailed derivation including the treatment of the oscillator bath. We emphasize that \( \kappa \) and \( \gamma \) are basis independent within our approximation.

2. Numerical approach

To treat Eq. (37) numerically we project it on the instantaneous eigenbasis \( |i \rangle \) of \( H_{\text{sys}} \) to get it on Redfield form

\[
\dot{\rho}_{ij} = -i \omega_{ij} \rho_{ij} - \sum_{kl} R_{ijkl} \rho_{kl},
\]

where \( \omega_{ij} = (E_i - E_j) / \hbar \) is the energy difference between the eigenstates \( |i \rangle \) and \( |j \rangle \) and \( \hat{R} \) denotes the Redfield tensor

\[
R_{ijkl} = \frac{\kappa}{2} \sum_{\mu=1}^{3} \left[ \delta_{ik} \delta_{lj} - i (|i| \sigma_\mu |k\rangle \langle l| \sigma_\mu^\dagger |j\rangle \langle j| - (i|k\rangle \langle l| \sigma_\mu |j\rangle \langle j| - (i|k\rangle \langle l| \sigma_\mu^\dagger |j\rangle \langle j| - (i|k\rangle \langle l| \sigma_\mu |j\rangle \langle j| - (i|k\rangle \langle l| \sigma_\mu^\dagger |j\rangle \langle j| - \delta_{ji} \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj} - 2 i \delta_{ij} \delta_{lj}.
\]

where \( \delta \) is the Kronecker delta and the summation \( \mu \) runs over the number of qubits. Numerically, it is convenient to work in the Liouvillian space where \( \rho \) is a vector quantity. In this way, we rewrite Eq. (38) as a matrix equation

\[
\dot{\rho}_{ij} + i \omega_{ij} \rho_{ij} = \sum_{kl} R_{ijkl} \rho_{kl}.
\]

The three qubits and oscillator span an \( (8N+1) \)-dimensional Hilbert space, where \( N \) is the number of photons in the resonator. This makes \( \rho \) a column vector of length \( 64(N+1)^2 \). The Redfield tensor is a \( 64(N+1)^2 \times 64(N+1)^2 \) matrix and \( \omega \) is a diagonal matrix of the same size. We work in the eigen-
basis of the qubit+oscillator Hamiltonian which we obtain by exact numerical diagonalization. The solution to Eq. (40) is given by

$$\rho(t) = e^{-(i\omega + \hat{R})t} \rho(0) = \hat{U}(t) \rho(0),$$  

which can be solved by numerically diagonalizing the propagator $\hat{U}(t)$. In this way we sequentially simulate the entire pulse sequence with the output density matrix of one passage serving as the input of the next. In the spirit of the previous sections all errors occurring during single-qubit operations are neglected. Hence, we are not concerned with the dynamics of these gates, which consequently are realized as matrix multiplications. As in the previous sections we use the minimum fidelity as the measure of code performance, which is obtained by numerically searching the space of initial states.

3. Numerical results

The goal of this section is to show that, for realistic values of qubit-oscillator coupling, the cavity-mediated cPhase gate outperforms the diagonal coupling in the limit $\gamma \rightarrow 0$ when the oscillator is not damped. This is very encouraging for future applications, since resonators with lifetimes several orders of magnitude longer than qubits can be made.$^{31,32}$

The gate operation time $t_g$ is now set by the qubit oscillator coupling and the level splitting of the qubits and oscillator, respectively. From experiments on similar systems we expect the coupling energy to be in the range $g/h \approx 10–100$ MHz (Ref. 2) for all qubits. To get restrictive results we chose $g/h=18$ MHz as a typical value. This corresponds to a temperature $T \sim 0.8$ mK, and is thus consistent with the high-temperature approximation discussed in Sec. IV A 1. The qubit dephasing rate was set to $\Gamma_c=1$ MHz in all simulations. The qubit energies where taken to be $E_{ij}/h$ =4850, 5000, and 5150 MHz, respectively. The energy separation was chosen to match the performance of current state of the art tunable oscillators in superconducting systems.$^{31,32}$

The oscillator frequencies in the idle periods were chosen to be $\omega_o/2\pi=4925$ and 5075 MHz, respectively. With these parameters the gate operation time is given by $\Gamma_c t_g=0.26,$$^{22}$ which is in the range where coding improves the fidelity for the diagonal coupling.

We plot the minimum fidelity normalized to the minimum no-coding fidelity for several values of the cavity dephasing rate $\gamma$ in Fig. 8. To compare the two implementations, the result for the diagonal gate [Eq. (20)] is plotted in the same figure (dotted line). When cavity and qubits are subject to the same amount of dephasing ($\gamma=\kappa$), the cavity-mediated coupling cannot match the diagonal. This is however to be expected since we have introduced an additional uncorrectable channel of noise through the cavity. It is however interesting to see that the fidelity can be improved with a modest reduction in cavity-environment coupling. With a choice of $\gamma =0.75\kappa$ the two couplings exhibit comparable fidelities as can be seen in Fig. 8. Further decrease in $\gamma$ improves the situation even further. If storage times up to $\Gamma_c t_g=2$ are considered, the cavity-mediated coupling actually performs better than the diagonal coupling. We attribute this improvement to the SWAP operations, which transfer the state from the qubit to the cavity. Given that the cavity is more phase coherent than the qubit the coherence is thus partially protected during part of the processing.

V. CONCLUSION

In this paper we have studied the performance of the three-qubit phase-flip QECC for realistic gates in superconducting systems. Since such a quantum code requires a limited amount of qubits for its implementation, it is interesting from the point of view of superconducting devices, where such an experiment should be realizable in the near future. We have studied two explicit couplings, one diagonal and one cavity mediated.

Our analysis begin with the case of ideal gates, where no gate errors occur during processing. In this case, we show that coding is beneficial, not only in the short-time limit, but can also significantly reduce the dephasing rate when considering realistic experimental parameters. We move on to consider realistic gates, deriving analytical expressions for the fidelity in the case where the qubits are coupled in diagonal fashion. We found an upper limit on the gate operation time $t_g$ allowed to benefit from coding. For the cavity-mediated coupling, we study the system numerically, solving the master equation that describes dephasing in the qubits and cavity. For realistic values of cavity-qubit coupling we find that the fidelity is comparable to that of the much simpler diagonal coupling. In the limit of weak coupling between the environment and cavity, this coupling even outperforms the diagonal one. We attribute this effect to the transfer of coherence from the qubit to the cavity, where it is protected during part of the pulse sequence. In view of the high $Q$ values demonstrated for stripline cavities, this is promising for future applications.
This work is supported by the European Commission through the IST-015708 EuroSQIP integrated project and by the Swedish Research Council.

APPENDIX A: DERIVATION OF THE MASTER EQUATION

The cavity plus qubit system with dephasing is modeled by coupling its constituents longitudinally to a thermal bath so that the Hamiltonian for the full system plus bath is given by Eq. (32). In the Markov approximation, it is possible to derive an equation of motion for the reduced cavity plus qubit density matrix. To second order in the couplings \( \lambda_i \), this equation reads

\[
\dot{\rho}_{ij} = -i \omega_{ij} \rho_{ij} + L_{ij}^{n,Y}(\rho) + L_{ij}^{e,X}(\rho),
\]

where \( \omega_{ij} = (E_i - E_j)/\hbar \) is the energy difference between the eigenstates \( |i\rangle \) and \( |j\rangle \) of \( H_{\text{sys}} \). The dissipative dynamics is governed by \( L(\rho) \) whose matrix elements are given by

\[
L_{ij}^{n,Y}(\rho) = \sum_{kl} \omega_{kl} \rho_{kl} L_0^Y(\omega_{kl}) - \omega_{ij} \rho_{ij} L_0^Y(\omega_{ij}),
\]

where \( \omega \) is the system operator that couples to the bath and \( L(\Omega) \) is the Laplace transform of the bath correlator

\[
L_0(\Omega) = \int_0^\infty dt e^{-\Omega t} \langle Q(t) Q(0) \rangle.
\]

For the case of a bath in thermal equilibrium, the real part of \( L_0(\Omega) \) is responsible for a small energy shift (Lamb shift) of the energy levels. Hence it will not be relevant for any pure phase damping, which is the central interest from the point of view of error correction. This effect will therefore be disregarded in the following analysis. The relevant transition frequencies \( \omega_{ij} \) will be determined by the matrix elements of \( \sigma_z \) and \( n \), respectively. We must therefore separate our analysis into two cases: the resonant and the off-resonant periods in the pulse sequence described in Sec. IV. In the off-resonant case, the eigenstates of the system are the product states \( |g/e; n\rangle \) for which both \( n \) and \( \sigma_z \) are diagonal. This leaves us with only one relevant transition frequency \( \omega_{ij} = 0 \). When the qubit and cavity are on resonance, the eigenstates of the system are the dressed states of the Jaynes-Cummings Hamiltonian

\[
|n; \pm\rangle = \frac{1}{\sqrt{2}}(|g; n\rangle \pm i|e; n - 1\rangle),
\]

for which the matrix elements of \( m \) and \( \sigma_z \) reads

\[
\langle m; \pm | \sigma_z | n; \pm \rangle = 0,
\]

\[
\langle m; \pm | n | n; \pm \rangle = \left( n - \frac{1}{2} \right) \delta_{mn},
\]

\[
\langle m; \pm | n | n; \mp \rangle = \frac{1}{2} \delta_{mn}.
\]

In this case we get three relevant frequencies \( \omega_{n,\pm,\mp} = \pm \frac{\gamma}{\hbar} \). We now assume that the baths have no structure on the scale of the qubit-cavity coupling. We further assume the temperature to be much higher than the cavity-qubit coupling \( k_B T \gg \hbar \gamma \). In this case we may safely approximate \( \text{Re} L_0(\pm \gamma) = \text{Re} L_0(0) \) and conclude that the only relevant parameter for dissipation will be \( L_0(0) \), for both the resonant and off-resonant regimes. With this clearly stated we can, from Eq. (A1), write down the master equation for the system dynamics as given in Eq. (37) with the rates \( \kappa \) and \( \gamma \) given by \( \kappa = 4 \text{Re} L_0(0) \) and \( \gamma = \text{Re} L_0(0) \) independently of the choice of \( H_{\text{sys}} \). These are the dephasing rates for the bare qubit and cavity systems, respectively.

29 C. W. Gardiner and P. Zoller, Quantum Noise (Springer-Verlag, Berlin, 2000).