Controlling Excitation Inversion of a Cooper Pair Box Interacting with a Nanomechanical Resonator *

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We investigate the action of time-dependent detunings upon the excitation inversion of a Cooper pair box interacting with a nanomechanical resonator. The method employs the Jaynes–Cummings model with damping, assuming different decay rates of the Cooper pair box and various fixed and time-dependent detunings. It is shown that, when the presence of damping plus constant detunings destroys the collapse/revival effects, convenient choices of time-dependent detunings allow one to reconstruct such events in a perfect way. It is also shown that the mean excitation of the nanomechanical resonator is more robust against damping of the Cooper pair box for convenient values of time-dependent detunings.

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A popular and exactly soluble model in quantum optics is the Jaynes–Cummings model (JCM). It describes the interaction of a two-level atom with a single-mode of the electromagnetic field.\textsuperscript{[1–8]} Over the last two decades various extensions of the ordinary JCM have been used in various contexts, e.g., (i) in the study of a three-level atom interacting with a two-mode squeezed vacuum,\textsuperscript{[9]} (ii) in the study of atom-field interaction inside a damped cavity,\textsuperscript{[10]} (iii) the same as in (i), including an additional (nonlinear) Kerr medium,\textsuperscript{[11,12]} (iv) two-level atoms inside a cavity acted upon by an external field control,\textsuperscript{[13]} (v) in the study of the nonlinear dynamical evolution of a driven two-photon JCM,\textsuperscript{[14]} (vi) a generalized JCM, including dissipation,\textsuperscript{[15–17]} including contexts of multiphoton interactions\textsuperscript{[30,19]} etc. In all these cases, with interest either in the field or in atomic properties the theoretical approach traditionally assumes the atom-field coupling as a constant parameter. Comparatively, the number of works in the literature is very small when such coupling and the atomic frequency are time-dependent parameters,\textsuperscript{[20–25]} including the case of time-dependent amplitudes.\textsuperscript{[26]} However, this scenario is also relevant; for example, the state of two qubits (qubits stand for quantum bits) with a desired degree of entanglement can be generated via a time-dependent atom-field coupling.\textsuperscript{[27]} Actually, such coupling can modify the dynamical properties of the atom and the field, with transitions that involve a large number of photons.\textsuperscript{[28]} In general, these studies are simplified by neglecting the atomic decay from an excited level. Theoretical treatments taking into account this complication of the real world may employ a modified JCM. As is expected, in this case the state describing the system decoheres, since the presence of dissipation destroys the state of a system as time flows.

In the present work we extend what we have learned from the applications of JCM to the atom-field interaction to investigate an advantageous system, from the experimental viewpoint (faster response, better controllability, and useful scalability for quantum computing).\textsuperscript{[29]} We consider a nanomechanical resonator (NR) interacting with a Cooper pair box (CPB).\textsuperscript{[16,18,40]} Such a nanosystem has been explored extensively in the recent literature, e.g., to investigate: (i) quantum nondemolition measurements,\textsuperscript{[31–32]} (ii) decoherence of nonclassical states, as Fock states and superposition states describing mesoscopic systems,\textsuperscript{[33–55]} etc. The fast advance in the technique of fabrication in nanotechnology has implied great interest in the study of the NR system in view of its potential applications, as a sensor, to be used in biology, astronomy, quantum computation, and quantum information\textsuperscript{[36–39]} to implement the quantum qubit\textsuperscript{[40]} and the generation of nonclassical states, as Fock state,\textsuperscript{[41–44]} Schrödinger’s cat state,\textsuperscript{[45]} squeezed states,\textsuperscript{[45]} clusters states\textsuperscript{[46]} etc. In particular, when accompanied by superconducting charge qubits, the NR has been used to prepare entangled states.\textsuperscript{[47]} Zhou et al.\textsuperscript{[45]} have proposed a scheme to prepare squeezed states using an NR coupled to a CPB qubit; in this proposal the NR-CPB coupling is under an external control whereas the connection between these two subsystems plays an important role in quantum computation. Such a control is achieved via a convenient change of the system parameters, which can set “on” and “off” the interaction between the NR and the CPB, on demand.

In this study we treat the CPB excitation inversion, its control, and the average photon number in the NR. We consider dissipation in the CPB due to a decay rate from an excited state to the ground state. We also verify in which way the time dependence of
the CPB-NR coupling modifies the mentioned properties. To this end we must solve the time evolution of the whole CPB-NR system, via the approach presented in the following.

\[ H = \Omega \hat{a}^\dagger \hat{a} + 4E_c(N_g - \frac{1}{2})\hat{\sigma}_z - 4E_0^0 \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) \hat{\sigma}_z, \]

where \( \hat{a} \) (\( \hat{a}^\dagger \)) is the creation (annihilation) operator for the NR excitation, with frequency \( \Omega \) and mass \( m \); \( E_0^0 \) and \( E_c \) are respectively the energy of each Josephson junction and the charge energy of a single electron; \( C_g \) and \( C_0^0 \) are the input capacitance and the capacitance of each Josephson tunnel, respectively. \( \Phi_0 = h/2e \) is the quantum flux and \( N_g = C_g V_1/2e \) is the charge number in the input with the input voltage \( V_1 \). We have used the Pauli matrices to describe our system operators, where the states \( |0\rangle \) and \( |1\rangle \) represent the number of extra Cooper pairs in the superconducting island. We have \( \hat{\sigma}_z = |1\rangle\langle 1| - |0\rangle\langle 0| \), \( \hat{\sigma}_x = |0\rangle\langle 1| + |1\rangle\langle 0| \) and \( E_C = e^2/(C_g + 4C_0^0) \).

The magnetic flux can be written as the sum of two terms,

\[ \Phi_e = \Phi_1 + B \ell \hat{x}, \]

where the first term \( \Phi_1 \) is the induced flux, corresponding to the equilibrium position of the NR and the second term describes the contribution due to the vibration of the NR; \( B \) represents the magnetic field created in the loop. We have assumed the displacement \( \hat{x} \) described as \( \hat{x} = x_0 (\hat{a}^\dagger + \hat{a}) \), where \( x_0 = \sqrt{\pi m \Omega/2} \) is the amplitude of the oscillation. Substituting Eq. (2) in Eq. (1) and controlling the flux \( \Phi_e \) we can adjust \( \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) = 0 \), making the approximation \( \pi B \ell x/\Phi_0 \ll 1 \) the above Hamiltonian results as (in rotating wave approximation),

\[ \hat{H} = \Omega \hat{a}^\dagger \hat{a} + \frac{1}{2} \omega_0 \hat{\sigma}_z + \lambda_0 (\hat{\sigma}_+ \hat{a} + \hat{a}^\dagger \hat{\sigma}_-), \]

where the constant coupling \( \lambda_0 = -4E_0^0 \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) \left( \pi B \ell x/\Phi_0 \right) \) and the effective energy \( \omega_0 = 8E_c(N_g - \frac{1}{2}) \). An important advantage of this coupling mechanism is its easy and convenient controllability.

Next, we extend the previous approach to a more general scenario by substituting \( \Omega(t) = \Omega + f(t) \) and \( \lambda_0 \rightarrow \lambda(t) = \lambda_0 [1 + f(t)/\Omega] \)\(^{[21,50,51]} \) in addition we assume the presence of a constant decay rate \( \gamma \) in the CPB; \( \omega_0 \) is the transition frequency of the CPB and \( \lambda_0 \) stands for the CPB-NR coupling. Here \( \hat{\sigma}_\pm \) and \( \hat{\sigma}_z \) are the CPB transition and excitation inversion operators, respectively; they act on the Hilbert space of atomic states and satisfy the commutation relations \( [\hat{\sigma}_+, \hat{\sigma}_-] = \hat{\sigma}_z \) and \( [\hat{\sigma}_z, \hat{\sigma}_\pm] = \pm 2\hat{\sigma}_\pm \). As is well known, the coupling parameter \( \lambda(t) \) is proportional to \( \sqrt{R(t)/V(t)} \), where the time-dependent quantization volume \( V(t) \) takes the form \( V(t) = V_0/\left[ 1 + f(t)/\Omega \right] \)\(^{[21,50,51]} \). Accordingly, we obtain the

Fig. 1. Model for the CPB-NM coupling.

A Josephson charge qubit system has been used to couple with an NR. Here we study a modified model where a CPB is coupled to an NR, as shown in Fig. 1. The scheme is inspired by the works of Liao et al.\(^{[48]} \) and Zhou et al.\(^{[49]} \) where we have substituted each Josephson junction by two of them. This creates a new configuration that includes a third loop. A superconducting CPB charge qubit is adjusted via a voltage \( V_1 \) at the system input and a capacitance \( C_g \). We require a scheme that attains an efficient tunneling effect for the Josephson energy. In Fig. 1 we can observe three loops: one great loop between two small ones. This makes it easier to control the external parameters of the system since the control mechanism includes the input voltage \( V_1 \) plus three external fluxes \( \Phi_L \), \( \Phi_r \) and \( \Phi_e \). In this way one can induce small neighboring loops. The great loop contains an NR which is modeled as a harmonic oscillator with a high-Q mode of frequency \( \Omega \) and its effective area in the center of the apparatus changes as the NR oscillates, which creates an external flux \( \Phi_e(t) \) that provides the CPB-NR coupling.

In pursuing the quantum behavior of a macro scale object the nano scale mechanical resonator plays an important role. At sufficiently low temperature the zero-point fluctuation of the NR will be comparable to its thermal Brownian motion. The detection of zero-point fluctuations of the NR can give a direct test of Heisenberg’s uncertainty principle. Assuming a sensitivity up to 10 times the amplitude of the zero-point fluctuation, LaHaye et al.\(^{[48]} \) have experimentally detected the vibrations of a 20 MHz mechanical beam of tens of micrometres size. For a 20 MHz mechanical resonator its temperature must be cooled below 1 mK to suppress the thermal fluctuation. For a GHz mechanical resonator a temperature of 50 mK is sufficient to effectively freeze out its thermal fluctuation and let it enter the quantum regime. This temperature is already attainable in dilution refrigerators.

In this work we assume the four Josephson junctions to be identical, with the same Josephson energy \( E_0^0 \), the same assumption for the external fluxes \( \Phi_L \) and \( \Phi_r \), i.e., with same magnitude but opposite sign: \( \Phi_L = -\Phi_r = \Phi_z \). This interaction actually couples the two subsystems. Together with the free Hamiltonian of flux qubit and NR, the Hamiltonian of the whole system reads

\[ \hat{H} = \Omega \hat{a}^\dagger \hat{a} + 4E_c(N_g - \frac{1}{2})\hat{\sigma}_z - 4E_0^0 \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) \cos \left( \frac{\pi \Phi_e}{\Phi_0} \right) \hat{\sigma}_z, \]
new (non hermitian) Hamiltonian
\[ \hat{H} = \Omega(t) \hat{a}^\dagger \hat{a} + \frac{1}{2} \omega_0 \hat{\sigma}_z + \lambda(t) (\hat{\sigma}_+ \hat{a} + \hat{\sigma}_- \hat{a}^\dagger) - i \frac{\gamma}{2} |1\rangle \langle 1|. \]

(4)

It is worth remembering that non Hermitian Hamiltonians (NHH) are largely used in the literature. To give a few examples we mention: Ref. [52], where the authors used an NHH and an algorithm to generalize the conventional theory; Ref. [53], using an NHH to get information about entrance and exit channels; Ref. [54], using non Hermitian techniques to study canonical transformations in quantum mechanics; Ref. [55], solving quantum master equations in terms of NHH; Ref. [56], using a new approach for NHH to study the spectral density of weak H-bonds involving damping; Ref. [57], studying NHH with real eigenvalues; Ref. [58], using a canonical formulation to study dissipative mechanics exhibiting complex eigenvalues; Ref. [59], studying NHH in non commutative space, and more recently: Ref. [60], studying the optical realization of relativistic NHH; Ref. [23], studying the evolution of entropy of atom-field interaction; Ref. [22], using a damping JC-Model to study entanglement between two atoms, each of them lying inside different cavities.

![Fig. 2. Time evolution of the excitation inversion in the CPB with (a) \( n = 25 \), \( \Omega = \omega_0 = 2000 \lambda_0 \), for \( f(t) = 0 \) (resonance) and different values of decay rates \( \gamma \): (a) \( \gamma = 0.01 \lambda_0 \), (b) \( \gamma = 0.05 \lambda_0 \) and (c) \( \gamma = 0.5 \lambda_0 \).](image)

To find the state that describes our time-dependent system we will write it in the form
\[ |\Psi(t)\rangle = \sum_{n=0}^{\infty} (C_{0,n}(t)|0,n\rangle + C_{1,n}(t)|1,n\rangle), \]

(5)

where \( |1,n\rangle \langle 0,n| \) represents the CPB in its excited state \( |1\rangle \) (ground state \( |0\rangle \)). We take the CPB initially prepared in its excited state \( |1\rangle \) and the NR in a coherent state \( |\alpha\rangle \). By expanding the coherent state in Fock’s basis we have \( |\alpha\rangle = \sum_{n=0}^{\infty} F_n(\alpha)|n\rangle \) with \( F_n(\alpha) = \exp(-|\alpha|^2/2)|\alpha|^n/\sqrt{n!} \). Assuming the NR and the CPB decoupled at \( t = 0 \) and the initial conditions \( C_{0,n}(0) = 0 \) and \( \sum_{n=0}^{\infty} |C_{1,n}(0)|^2 = 1 \) we can write Eq. (5) as \( |\Psi(0)\rangle = \sum_{n=0}^{\infty} F_n(\alpha)|1,n\rangle \).

The time-dependent Schrödinger equation for this system is
\[ i \frac{d|\Psi(t)\rangle}{dt} = \hat{H} |\Psi(t)\rangle, \]

(6)

with the Hamiltonian \( \hat{H} \) given in Eq. (4). Substituting Eq. (4) in Eq. (6) we obtain the (coupled) equations of motion for the probability amplitudes \( C_{1,n}(t) \) and \( C_{0,n+1}(t) \):
\[ \frac{\partial C_{1,n}(t)}{\partial t} = -i n \Omega(t) C_{1,n}(t) - i \frac{\gamma}{2} C_{1,n}(t), \]
\[ \frac{\partial C_{0,n+1}(t)}{\partial t} = -i (n+1) \Omega(t) C_{0,n+1}(t) - i \frac{\gamma}{2} C_{0,n+1}(t). \]

(7)

(8)

The numerical solutions for the coefficients \( C_{1,n}(t) \) and \( C_{0,n+1}(t) \) furnish the quantum dynamical properties of the system, including the CPB-NR entanglement.

For the cases \( f(t) = 0 \) and \( f(t) = \text{const.} \), Eqs. (7) and (8) are exactly soluble. We find, analytically,
\[ C_{0,n+1}(t) = -4i \lambda F_n \sqrt{n} - 1^ {1/43t} \sinh(1/4\zeta t)/\zeta, \]
\[ C_{1,n}(t) = 2\xi F_n (i \gamma + 2\Omega - 2\omega_0) \sinh(1/4\zeta t)/\zeta, \]
\[ - i \xi \cosh(1/4\zeta t), \]

(9)

where \( \zeta = [\gamma (\gamma + 4\Omega (\omega_0 - \Omega)) - 4(\Omega^2 + \omega_0^2) - 16\lambda^2(1 + n) + 8\Omega \omega_0]^{1/2}, \delta = \gamma + 2\Omega (1 + 2n) \) and \( \xi = i(1/2\zeta) e^{-1/43t}. \) However, when the coupling \( f(t) \) is time-dependent the solution to Eqs. (7) and (8) is found only numerically.

As is well known, in the presence of a decay rate \( \gamma \) in the CPB the state of the entire CPB-NR system becomes mixed. In this case its description requires the use of the density operator \( \hat{\rho}_{\text{NR}} \), and to describe the CPB (NR) subsystem we must trace over the variables of the NR (CPB) subsystem. This allows us to get, e.g., the reduced operator \( \hat{\rho}_{\text{NR}} = \text{Tr}_{\text{CPB}}(\hat{\rho}_{\text{CN}}) \) concerning the NR subsystem, namely,
\[ \hat{\rho}_{\text{NR}} = \sum_{n,n'=0}^{\infty} [C_{1,n}(t)C_{1,n'}^*(t) + C_{0,n}(t)C_{0,n'}^*(t)]|n\rangle\langle n'|. \]

(10)

The CPB excitation inversion, here denoted as \( I(t)_{\text{CPB}} \), is an important observable of two level systems. It is defined as the difference of probabilities of finding the system in the excited and ground state, as follows:
\[ I(t)_{\text{CPB}} = \sum_{n=0}^{\infty} |C_{1,n}(t)|^2 - |C_{0,n+1}(t)|^2. \]

(11)

Equation (12) allows one to track the time evolution of the CPB excitation inversion. Firstly, we assume the resonant case \( f(t) = 0 \) for different values of the decay rate \( \gamma \), with \( \alpha = 5 \) and \( \Omega = \omega_0 = 2000 \lambda_0 \) and take the NR initially in a coherent state with the average number of excitations \( \langle n \rangle = 25 \) as shown in Fig. 2. With exceptions of amplitudes, the plots
in Figs. 2(a), 2(b) and 2(c) are identical for collapse-revival effects; as expected, the higher the decay rate, the lower the amplitude of oscillation. In the presence of detuning, with \( f(t) = \Delta = \text{const} \), \( \Delta \ll \omega_0 \), \( \Omega \) we can see that the excitation inversion, displayed in Fig. 3(a), occurs within the interval \( 30 < \lambda_0t < 50 \) whereas in Fig. 3(b) it occurs in the range \( 60 < \lambda_0t < 75 \) and in Fig. 3(c) this effect rapidly vanishes.

Next, considering the case of a variable detuning as \( f(t) = \eta \sin(\omega t) \), we can see in the plots in Figs. 4(a), 4(b) and 4(c) that excitation inversion occurs frequently in Fig. 4(a) and it disappears when the parameter \( \eta \) increases, as shown in the plots (b) and (c). Now, even considering the worst results obtained in the off-resonant cases, with detuning \( \Delta = \eta = 60\lambda_0 \) as shown in Fig. 3(c) and Fig. 4(c) we can see that the collapse-revival effects are recovered via the increase of the parameter \( \omega' \), as shown in the plots in Figs. 5(a), 5(b) and 5(c). Thus, the parameter \( \omega' \) plays an important role in the control of collapse and revival effect in the CPB excitation inversion.

Figure 6 shows the NR average excitations in the presence of CPB decay rate for various values of amplitude of oscillations (parameter \( \eta \)). Plots (d) and (h) concern the resonant case: in (d) the decay rate is greater than that in (h). Plots (a), (b), and (c) are for constant decay rates, with (time independent) detuning that increases from (c) \( \rightarrow \) (b) \( \rightarrow \) (a). Finally, plots (e), (f) and (g) are for time-dependent detuning, with the parameter \( \eta \) increasing from (e) \( \rightarrow \) (f) \( \rightarrow \) (g). We note that the three plots for time-dependent detunings (e), (f) and (g) are better than those for constant detunings (a), (b) and (c); despite all plots being concerned with the same decay rate, the first group is more robust against decay. For example, comparing the plots (a) and (g): although in (a) the fixed detuning is \( \Delta = 60\lambda_0 \) and in (g) the maximum detuning is
\[ \Delta_{\text{max}} = \eta = 60\lambda_0 \] we can see that in the last case the decay of the NR average excitation is softer. One also observes from Fig. 5(c) that the interval \( \lambda_{0\text{t}} \), where the time-dependent detuning recovers the collapse-revival effect, coincides with that in Fig. 6(g) where the average excitation is around 5 times greater than that for constant detuning (Fig. 6(a)).

In summary, we have considered a Hamiltonian model that describes a CPB-NR interacting system to study the CPB excitation inversion, \( I(t)_{\text{CPB}} \), and the average excitation number of the NR, \( \langle n(t) \rangle_{\text{NR}} \). We have also considered the off-resonant case, with various values of the detuning parameter \( f = f_0: f = \Delta \); and \( f = \eta\sin(\omega't) \) and in the presence of CPB decay [about 10 times greater than the (then neglected) NR decay]. These properties are characteristics of the entangled state that describe this coupled system for various values of the parameters involved. We have assumed the CPB initially in its excited state and the NR initially in a coherent state (see preparation in Ref. [30]). Thus, the following three scenarios were treated: (i) both subsystems in resonance (detuning \( f = 0 \)); (ii) off-resonance, with a constant detuning \( f = \Delta \neq 0 \), and (iii) with a time-dependent (sineoidal) detuning \( f(t) = \eta\sin(\omega't) \). The results were discussed above: in summary, concerning the CPB excitation inversion, an interesting result emerged: although the presence of a constant detuning destroys the collapse and revivals of the excitation inversion, these effects are restituted by the action of convenient time-dependent detunings, even in the presence of the CPB damping. Concerning the NR average excitation number, another interesting result appears: convenient choices of the time-dependent detuning \( f(t) \) makes the NR subsystem more robust against the decay affecting the CPB subsystem. For constant values of detuning, our numerical results are similar to others in the literature using a master equation (see, e.g., Ref. [16]). Finally we emphasize that the change in magnetic flux \( \Phi_e \) (cf. Fig. 1), due to the presence of an external force upon the NR, is responsible for controlling the parameters \( \Omega(t) \) and \( \lambda(t) \). From a future perspective, the study of the influence of other types of time-dependent detunings on these and other properties deserves attention, e.g., the influence of such detunings on the nonclassical depth of a subsystem state.

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