Quantum communication via controlled holes in the statistical distribution of excitations in a nanoresonator coupled to a Cooper pair box*

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We propose a scheme to transmit information via the statistical distribution of excitations of a nanomechanical resonator. It employs a controllable coupling between this system and a Cooper pair box. The success probability and the fidelity are calculated and compared with those obtained in an atom–field system in different regimes. Addtionaly, the scheme can also be applied to prepare low excited Fock states.

Keywords: quantum communication, quantum state engineering, superconducting circuits, nanomechanical resonator

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1. Introduction

Nanomechanical resonators (NRs) have been studied in a diverse number of situations, such as weak force detections,[1] precision measurements,[2] and quantum information processing.[3] The demonstration of the quantum nature of mechanical and micromechanical devices is a pursuing target, for example, the manifestations of purely nonclassical behavior in a linear resonator should exhibit energy quantization, the appearance of Fock states, quantum limited position–momentum uncertainty, superposition and entangled states, etc. The NRs can now be fabricated with fundamental vibrational mode frequencies in the MHz–GHz range.[4−6] Advances in the research of micromechanical devices also raise a fundamental question, whether such systems that contain macroscopic numbers of atoms will exhibit the quantum behavior. Due to their sizes, the quantum behavior in micromechanical systems will be strongly influenced by their interactions with the environment, and the existence of an experimentally accessible quantum regime will depend on the rate at which decoherence occurs.[7,8] One crucial step in the study of nanomechanical systems is the engineering and detection of the quantum effects of the mechanical modes.[9] This can be achieved by connecting the resonators with solid-state electronic devices,[11−14] such as a single-electron transistor. The NR has also been used to study quantum nondemolition measurement,[14−17] quantum decoherence,[13,18] and macroscopic quantum coherence phenomena.[19] The fast advance in the nanotechnology raises great interest in the study of the NR system due to its potential modern applications, such as sensors largely used in various fields including biology, astronomy, and quantum computation.[20,21] It can also be used to implement the quantum qubit[22−24] and the quantum multiqubit[25] in quantum information[3,22,26−30] and to explore cooling mechanisms,[31−39] transducer techniques,[40−42] and the generation of nonclassical states (such as the Fock[43] state, Schrödinger's cat,[13,44,45] the squeezed states[46−50] including intermediate and other superposition states.[51−55]) And more recently, it also find applications in atomic physics and quantum optics.[56] In particular, the NR coupled with superconducting charge qubits has been used to generate entangled states.[13,44,57,58] In a previous paper, Zhou and Mizel[49] proposed a scheme to create squeezed states in an NR coupled to a Cooper pair box (CPB) qubit, where the NR–CPB coupling

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was controllable. Sucha control comes from the change of the external parameters and plays an important role in quantum computation, allowing us to turn on and off the interaction between the systems on demand.

Now, the storage of optical data and communications using basic processes of quantum physics have been subjects of growing interest in recent years.[59] We present here a feasible experimental scheme to create holes in the statistical distribution of excitations of a coherent state previously prepared in an NR. In this proposal, the coupling strength between the NR and the CPB can be controlled continuously by tuning two external biasing fluxes. The motivation is inspired by the early investigations on the production of new materials possessing holes in their fluorescent spectra[60] and is also inspired by our previous work, in which we have used alternative systems and schemes to attain the goal.[61−64] Producing holes with controlled positions in the number space is desirable due to their possible applications in quantum computation, quantum cryptography, and quantum communication. As argued in Ref. [63], these states may be used for optical data storage, each hole being associated with some signals (say yes, $|1\rangle$, or $|+\rangle$) and its absence being associated with the opposite signals (no, $|0\rangle$, or $|-\rangle$). The generation of such holes has been discussed in the contexts of cavity-QED[64] and traveling waves.[65]

2. Model for a CPB–NR system

There exist in the literature a large number of devices using the SQUID, where the CPB charge qubit consists of two superconducting Josephson junctions in a loop. In the present model, a CPB is coupled to an NR, as shown in Fig. 1. The scheme is inspired by the works by Liao *et al.*[27] and Zhou *et al.*,[49] and we have substituted each Josephson junction by two of them. This creates a new configuration including a third loop. A superconducting CPB charge qubit is adjusted via a voltage V_1 at the system input and a capacitance C_1Josephson energy. In Fig. 1, we can see three loops, one great loop connecting two small ones. This makes it easier to control the external parameters of the system, since the control parameters include the input voltage V_1 plus three external fluxes $\Phi_{\rm L}$, $\Phi_{\rm r}$, and $\Phi_{\rm e}(t)$. In this way, we can induce small neighboring loops. The great loop contains the NR, and its effective area in the center of the apparatus changes as the NR oscillates, which creates an external flux $\Phi_{\rm e}(t)$ that provides the CPB–NR coupling to the system.

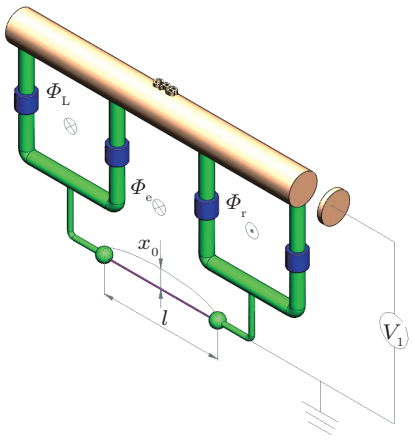


Fig. 1. (colour online) Model for the CPB–NR coupling.

In this work, we assume the four Josephson junctions being identical with the same Josephson energy $E_{\rm J}^0$. The external fluxes $\Phi_{\rm L}$ and $\Phi_{\rm r}$ are assumed to have the same magnitude but opposite signs, $\Phi_{\rm L} = -\Phi_{\rm r} = \Phi_x$. In this way, we can write the Hamiltonian describing the entire system as

$$\hat{H} = \omega\hat{a}^{\dagger}\hat{a} + 4E_{\rm c}\left(N_1 - \frac{1}{2}\right)\hat{\sigma}_z - 4E_{\rm J}^0\cos\left(\frac{\pi\Phi_x}{\Phi_0}\right)\cos\left(\frac{\pi\Phi_{\rm e}}{\Phi_0}\right)\hat{\sigma}_x, \quad (1)$$

where $\hat{a}^{\dagger}$ ($\hat{a}$) is the creation (annihilation) operator for the NR excitation with frequency ω and mass m; $E_{\rm c}$ is the charge energy of a single electron; C_1 and $C_{\rm J}^0$ stand for the input capacitance and the capacitance of the Josephson tunnel, respectively; $\Phi_0 = h/2e$ is the quantum flux; and $N_1 = C_1V_1/2e$ 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 states $|g\rangle$ and $|e\rangle$ (or 0 and 1) represent the numbers of extra Cooper pairs in the superconduting island. We have $\hat{\sigma}_z = |g\rangle\langle g| - |e\rangle\langle e|$, $\hat{\sigma}_x = |g\rangle\langle e| - |e\rangle\langle g|$, and $E_{\rm c} = e^2/(C_1 + 4C_{\rm J}^0)$.

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

$$\Phi_{\rm e} = \Phi_{\rm b} + Bl\hat{x}, \quad (2)$$

where $\Phi_{\rm b}$ 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, with B being the magnetic field created in the loop. We have assumed that the displacement $\hat{x}$ is described as $\hat{x} = x_0(\hat{a}^{\dagger} + \hat{a})$, where $x_0 = \sqrt{m\omega/2}$ is the amplitude of the oscillation.

Substituting Eq. (2) in Eq. (1) and controlling flux $\Phi_{\rm b}$, we can adjust $\cos(\pi\Phi_{\rm b}/\Phi_0)=0$ to obtain

$$\hat{H} = \omega\hat{a}^{\dagger}\hat{a} + 4E_{\rm c}\left(N_1-\frac{1}{2}\right)\hat{\sigma}_z - 4E_{\rm J}^0\cos\left(\frac{\pi\Phi_x}{\Phi_0}\right)\sin\left(\frac{\pi Bl\hat{x}}{\Phi_0}\right)\hat{\sigma}_x. \quad (3)$$

Making the approximation $\pi Blx/\Phi_0 \ll 1$, we find

$$\hat{H} = \omega\hat{a}^{\dagger}\hat{a} + \frac{1}{2}\omega_0\hat{\sigma}_z + \lambda_0(\hat{a}^{\dagger}+\hat{a})\hat{\sigma}_x, \quad (4)$$

where the constant coupling is

$$\lambda_0 = -4E_{\rm J}^0\cos\left(\frac{\pi\Phi_x}{\Phi_0}\right)\left(\frac{\pi Blx_0}{\Phi_0}\right),$$

and the effective energy is $\omega_0 = 8E_{\rm c}(N_1-\frac{1}{2})$. Under the rotating wave approximation, the above Hamiltonian becomes

$$\hat{H} = \omega\hat{a}^{\dagger}\hat{a} + \frac{1}{2}\omega_0\hat{\sigma}_z + \lambda_0(\hat{\sigma}_+\hat{a} + \hat{a}^{\dagger}\hat{\sigma}_-). \quad (5)$$

In the interaction picture, the Hamiltonian is written as

$$\hat{H}_{\rm I} = \hat{U}_0^{\dagger}\hat{H}\hat{U}_0 - {\rm i}\,\hbar\hat{U}_0^{\dagger}\frac{\partial\hat{U}_0}{\partial t},$$

where

$$\hat{U}_0 = \exp\left[-{\rm i}\left(\omega\hat{a}^{\dagger}\hat{a} + \frac{\omega_0\hat{\sigma}_z}{2}\right)t\right]$$

is the evolution operator in the absence of the CPB–NR interaction. Assuming the system operating under the resonant condition, i.e., $\omega=\omega_0$, and setting $\hat{\sigma}_z = \hat{\sigma}_+\hat{\sigma}_- - \hat{\sigma}_-\hat{\sigma}_+$ and $\hat{\sigma}_\pm = (\hat{\sigma}_x \pm {\rm i}\,\hat{\sigma}_y)/2$, with $\hat{\sigma}_y = (|e\rangle\langle g| - |e\rangle\langle g|)/{\rm i}$, we have the abbreviated form for the interaction Hamiltonian

$$\hat{H}_{\rm I} = \beta(\hat{a}^{\dagger}\hat{\sigma}_- + \hat{a}\hat{\sigma}_+), \quad (6)$$

where $\beta=-\lambda_0$, and $\hat{\sigma}_+ = |g\rangle\langle e|$ ($\hat{\sigma}_- = |e\rangle\langle g|$) is the raising (lowering) operator for the CPB.

We note that coupling constant β can be controlled through flux Φ_x, which influences the small loops in Fig. 1. Furthermore, we can control gate charge N_1 via the gate voltage V_1 syntonized to the coupling. It should be mentioned that energy ω_0 depends on the induced flux Φ_x. So, when we syntonize the induced flux Φ_x, energy ω_0 changes. To avoid unnecessary transitions, we assume the changes in the flux are slow enough to obey the adiabatic condition.

Next we show how to make holes in the statistical distribution of excitations in the NR. We start with the CPB initially prepared in its ground state $|{\rm CPB}\rangle = |g\rangle$ and the NR initially prepared in the coherent state $|{\rm NR}\rangle = |\alpha\rangle$. Then state $|\Psi\rangle$ that describes the entire CPB–NR system evolves as

$$|\Psi_{\rm NC}(t)\rangle = \hat{U}(t)|g\rangle|\alpha\rangle, \quad (7)$$

where $\hat{U}(t) = \exp(-{\rm i}\,t\hat{H}_{\rm I})$ is the evolution operator in the interaction picture. Then after some algebra, we obtain

$$\hat{U}(t) = \cos(\beta t\sqrt{\hat{a}^{\dagger}\hat{a}+1})|g\rangle\langle g| + \cos(\beta t\sqrt{\hat{a}^{\dagger}\hat{a}})|e\rangle\langle e| - {\rm i}\,\frac{\sin(\beta t\sqrt{\hat{a}^{\dagger}\hat{a}+1})}{\sqrt{\hat{a}^{\dagger}\hat{a}+1}}\hat{a}|g\rangle\langle e| - {\rm i}\,\frac{\sin(\beta t\sqrt{\hat{a}^{\dagger}\hat{a}})}{\sqrt{\hat{a}^{\dagger}\hat{a}}}\hat{a}^{\dagger}|e\rangle\langle g|. \quad (8)$$

In this way, the evolved state in Eq. (7) becomes

$$|\Psi_{\rm NC}(t)\rangle = {\rm e}^{-\frac{|\alpha|^2}{2}}\sum_{n=0}^{\infty}\frac{\alpha^n}{\sqrt{n!}}[\cos(\omega_n\tau)|g,n\rangle - {\rm i}\sin(\omega_n\tau)|e,n+1\rangle], \quad (9)$$

where $\omega_n = \beta\sqrt{n+1}$. If we detect the CPB in state $|g\rangle$ after a convenient time interval τ_1, then state $|\Psi_{\rm NC}(t)\rangle$ reads

$$|\Psi_{\rm NC}(\tau_1)\rangle = \eta_1\sum_{n=0}^{\infty}\frac{\alpha^n}{\sqrt{n!}}\cos(\omega_n\tau_1)|n\rangle, \quad (10)$$

where η_1 is a normalization factor. If we choose τ_1 in such a way that $\beta\sqrt{n_1+1}\tau_1=\pi/2$, then component $|n_1\rangle$ in Eq. (10) is eliminated.

In the second step, we assume that the first CPB is rapidly replaced by a second one in state $|g\rangle$, which interacts with the NR after the above detection (instead of replacing the CPB by another one, we can also repeat successively the operation with the same CPB). For the second CPB, the initial state of the NR is the state given in Eq. (10) produced by the detection of the first CPB in $|g\rangle$. As a result, the new CPB–NR system evolves to the state

$$|\Psi_{\rm NC}(\tau_2)\rangle = \sum_{n=0}^{\infty}\frac{\alpha^n}{\sqrt{n!}}[\cos(\omega_n\tau_2)\cos(\omega_n\tau_1)|g,n\rangle - {\rm i}\cos(\omega_n\tau_1)\sin(\omega_n\tau_2)|e,n+1\rangle]. \quad (11)$$

The detection of the second CPB again in state $|g\rangle$ leads the entire system to collapse to the state

$$|\Psi_{\rm NC}(\tau_2)\rangle = \eta_2\sum_{n=0}^{\infty}\frac{\alpha^n}{\sqrt{n!}}[\cos(\omega_n\tau_2)\cos(\omega_n\tau_1)|n\rangle], \quad (12)$$

where η_2 is a normalization factor. In this way, the choice $\beta\sqrt{n_2+1}\tau_2=\pi/2$ makes a second hole, now in component $|n_2\rangle$.

By repeating the above procedure M times, we obtain the generalized result for the M-th CPB detection

$$|\Psi_{\rm NC}(\tau_M)\rangle = \eta_M \sum_{n=0}^{\infty} \frac{\alpha^n}{\sqrt{n!}} \prod_{j=1}^{M} \cos(\omega_n \tau_j)|n\rangle, \quad (13)$$

where τ_j is the time spent in the j-th NR–CPB interaction. According to Eq. (13), the number of CPB states detected coincides with the number of holes produced in the statistical distribution. In fact, equation (13) allows us to find the expression for statistical distribution $P_n = |\langle n|\Psi_{\rm NC}(\tau_M)|^2$. A little algebra furnishes

$$P_n = \frac{(\alpha^{2n}/n!)\prod_{j=1}^{M}\cos^2(\omega_n\tau_j)}{\sum_{m=0}^{\infty}(\alpha^{2m}/m!)\prod_{j=1}^{M}\cos^2(\omega_n\tau_j)}\,. \quad (14)$$

To illustrate the results, we show the controlled production of holes in the photon number distribution in Fig. 2.

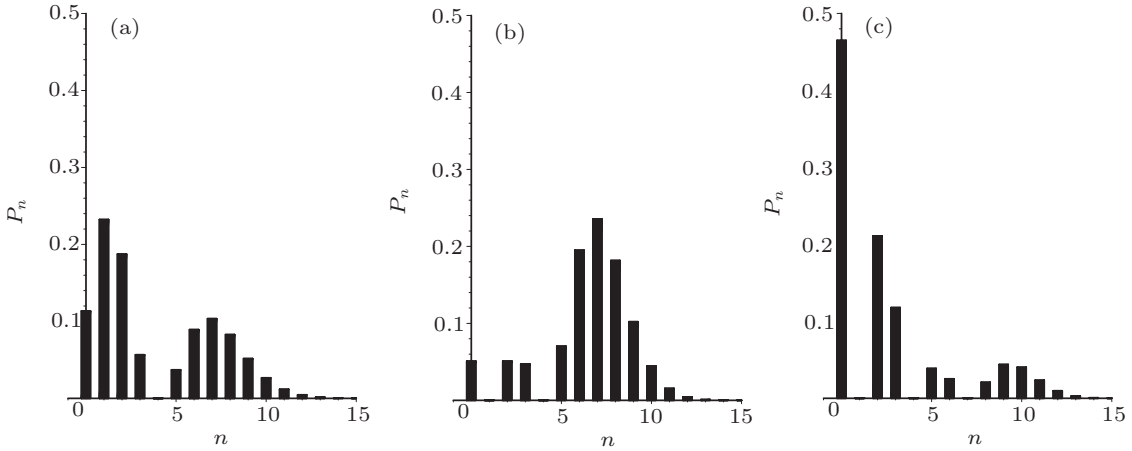


Fig. 2. Holes in the photon number distribution with $\alpha = 2.0$ (a) at $n_1 = 4$ for the 1-st step, (b) at $n_1 = 4$ and $n_2 = 1$ for the 2-nd step, and (c) at $n_1 = 4$, $n_2 = 1$, and $n_3 = 7$ for the 3-rd step.

The success probability to produce the desired state is given by

$$P_{\rm s} = {\rm e}^{-|\alpha|^2} \sum_{m=0}^{\infty} (\alpha^{2m}/m!) \prod_{j=1}^{M} \cos^2(\omega_n \tau_j). \quad (15)$$

Note that the holes shown in Figs. 2(a), 2(b), and 2(c) occur with success probabilities of 9%, 4%, and 0.3%, respectively.

We can take advantage of the this procedure and apply it to the engineering of nonclassical states, e.g., to prepare the Fock states[66] and their superposition states.[67] To this end, two strategies can be used. In the first one, we eliminate the components on the left and the right sides of a desired Fock state $|N\rangle$, namely, $|N-1\rangle$, $|N-2\rangle, \ldots$ and $|N+1\rangle, |N+2\rangle, \ldots$ In the second one, we only eliminate the left side components of the Fock state $|N\rangle$. In both cases, it is convenient to consider the final state of the NR as

$$|\Psi_{\rm NC}(\tau_M)\rangle' = \eta'_M \sum_{n=0}^{\infty} \frac{\alpha^n}{\sqrt{n!}} (-{\rm i})^M \times \prod_{j=1}^{M} \sin(\omega_{n+j}\tau_j)|n+M\rangle, \quad (16)$$

which can be easily obtained by detecting the Cooper pair box in state $|e\rangle$. The success probability $P'_{\rm s}$ to produce a Fock state $|N\rangle$ reads

$$P'_{\rm s} = {\rm e}^{-|\alpha|^2} \sum_{m=0}^{\infty} (\alpha^{2m}/m!) \prod_{j=1}^{M} \sin^2(\omega_{n+j}\tau_j). \quad (17)$$

In the first strategy, we prepare Fock states $|N\rangle$ with $N = M$, i.e., the phonon number N coincides with the number of CPB detections M. The fidelity of these states is given by the phonon number distribution at P_M associated with state $|\Psi_{\rm NC}(\tau_M)\rangle'$

$$P_M = \frac{\prod_{j=1}^{M}\sin^2(\sqrt{j}\beta\tau_j)}{\sum_{n=0}^{\infty}(\alpha^{2n}/n!)\prod_{n=1}^{M}\sin^2(\sqrt{n+j}\beta\tau_n)}. \quad (18)$$

We note that in this case the fidelity coincides with the N-th component of the statistical distribution P_n. Figure 3 shows the phonon-number distributions exhibiting the creation of Fock states $|3\rangle$, $|4\rangle$, and $|5\rangle$ all with a fidelity of 99%. And the NR in the initial coherent state has $\alpha = 0.6$.

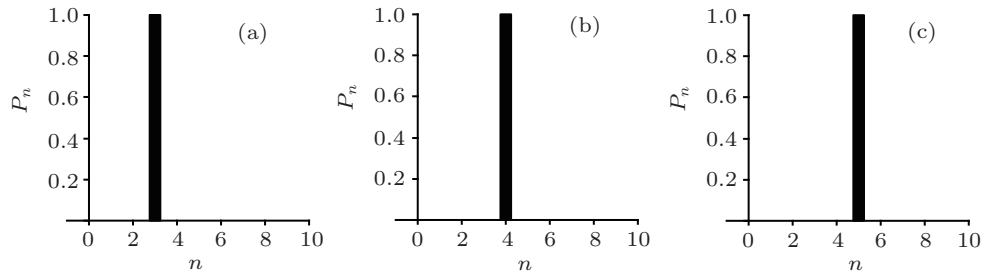


Fig. 3. Phonon number distributions exhibiting the creation of Fock states (a) $|3\rangle$ ($P'_{\rm s} = 17\%$), (b) $|4\rangle$ ($P'_{\rm s} = 11\%$), and (c) $|5\rangle$ ($P'_{\rm s} = 7\%$) all with a fidelity of 99%. The $\alpha = 0.6$ in the initial coherent state.

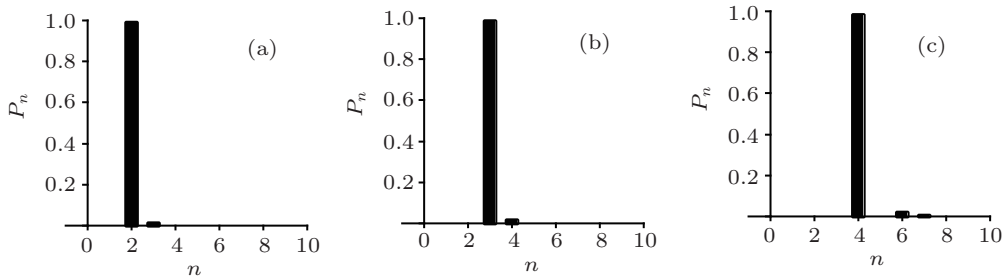


Fig. 4. Phonon number distributions exhibiting the creation of Fock states (a) $|2\rangle$ ($P'_{\rm s} = 17\%$), (b) $|3\rangle$ ($P'_{\rm s} = 1\%$), and (c) $|4\rangle$ ($P'_{\rm s} = 0.3\%$) all with a fidelity of 99%. The $\alpha = 0.6$ in the initial coherent state.

In the second strategy, we prepare Fock states $|N\rangle$ with $N = 2M$ or $2M - 1$. The associated fidelity is also given by Eq. (18). Figure 4 shows the phonon-number distributions exhibiting the creation of Fock states $|3\rangle$, $|4\rangle$, and $|5\rangle$ all with the same fidelity of 99%. And the NR in the initial coherent state has $\alpha = 0.6$.

3. Conclusion

For the feasibility of the scheme, it is worth mentioning some experimental values of parameters and characteristics of our system, the maximum value of the coupling constant is $\beta_{\max} \approx 45$ MHz, with $B \approx 0.1$ T, $l = 30$ μm, $x_0 = 500$ fm, $E_{\rm J}^0 = 5$ GHz, and $\omega_0 = 200\pi$ MHz.[6,27,48,49,68−72] The expression choosing the time spent to make a hole, $\beta\sqrt{n_j + 1}\tau_j = \pi/2$, gives $\tau_j \simeq 0.3$ ns by assuming all the CPB previously prepared at $t = 0$. On the other hand, the decoherence times of the CPB and the NR are 500 ns and 160 μs, respectively.[71] Accordingly, we may create about 1600 holes before the destructive action of the decoherence. However, when considering the success probability to detect all CPB in state $|g\rangle$, a more realistic estimation drastically reduces the number of holes. A similar situation occurs in Refs. [62]–[64] using the atom–field system to make holes in the statistical distribution P_n of a field state. In that case, about 1 μs is spent to create a hole, whereas 1 ms is the decoherence time of the field state inside the cavity. So, compared to that scenario, the present system is about 60% more efficient. For the generation of Fock state $|N\rangle$, equation (18) shows the convenience of starting with a low excited coherent state, since it involves only a small number of Fock components to be deleted via this hole burning procedure. The success probability is drastically reduced when we have to delete many components of the initial state to get the Fock state $|N\rangle$. As a consequence, this method will work only for small values of N ($N \lesssim 5$).

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