A single photon transistor based on superconducting systems

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We present a realistic scheme for how to construct a single-photon transistor where the presence or absence of a single microwave photon controls the propagation of a subsequent strong signal field. The proposal is designed to work with existing superconducting artificial atoms coupled to cavities. We study analytically and numerically the efficiency and the gain of our proposal and show that current transmon qubits allow for error probabilities \( \sim 1\% \) and gains of the order of hundreds.

In analogy with electronic transistors, a single photon transistor is a device where the presence or absence of a single gate photon controls the propagation of a large number of signal photons [1,2]. Such devices would represent a milestone enabling a plethora of new approaches for processing light, but their realization is hampered by the absence of interactions between photons. A promising route towards strong interactions at the single photon level consists of coupling propagating photons to individual atom-like systems [3,4]. The best realization of such a coupling is achieved in the microwave regime where experiments have demonstrated an unprecedented control of the coupling between superconducting artificial atoms and microwave photons [10,11]. In this letter we describe how to realize a single photon transistor based on existing superconducting technology. The resulting devices can be directly employed to detect individual itinerant microwave photons, and may find a range of applications within quantum information processing.

Various protocols for single photon non-linearities, quantum gates, and transistors using superconducting systems have been proposed [12-15], but have relied on unconventional qubit designs, multiple qubits and isolators, or temporal switching of parameters. A different method to construct the desired single photon transistors was described in Ref. [1], and realized in the optical regime in Ref. [2]. The atomic level structure used in that work is, however, not compatible with current superconducting artificial atoms (for convenience these will be referred to as atoms below). Here we show how to realize this proposal for a three level ladder (Ξ) system coupled to a superconducting cavity. This is the generic level scheme for the so-called transmon [16,20] and phase qubits [21] which essentially constitutes anharmonic ladders.

We focus on the three lowest levels of the system. We assume that the Ξ-system is sufficiently anharmonic such that the cavity is resonant with the upper transition between states \(|g\rangle\) and \(|e\rangle\), which has a coupling constant \(g\), but off-resonant from the lower transition between \(|f\rangle\) and \(|g\rangle\), see Fig. 1(a). For now, we ignore the cavity coupling on the lower transition, deferring the discussion of this off-resonant cavity coupling to later. With this arrangement, we can be in a regime where there is a strong cavity enhancement of the decay of the upper level \(|e\rangle\) via the Purcell effect, such that it decays rapidly compared to the two lower levels \(|g\rangle\) and \(|f\rangle\). In essence, this combined system of the Ξ-atom and the cavity thus has two metastable levels \(|f\rangle\) and \(|g\rangle\) which mimic the stable states of the A-type atom in the proposal of Ref. [1]. To control the system the transition between \(|f\rangle\) and \(|g\rangle\) is driven by a classical field. Below we describe how these ingredients allow us to implement a single photon transistor. For simplicity we discuss the scheme for a two sided cavity with identical (photon) decay rates \(\kappa\) to modes propagating to the left and right. Alternatively, the scheme could be realized with a single sided cavity if additional linear optics elements are used [22].

We describe the system using the “quantum jump” approach [30], where the evolution is described by a non-Hermitian Hamiltonian interrupted by quantum jumps at random times. The non-Hermitian Hamiltonian is given by

\[
H = \sum_{m=f,g,e} h(\omega_m - i\gamma_m/2)\hat{\sigma}_{mm} + h\omega_c\hat{a}^{\dagger}\hat{a}
+ \sum_{m=L,R} h \int dk \, c(k)\hat{b}_m^{\dagger}(k)\hat{b}_m(k) + i\hbar g [\hat{a}^{\dagger}\hat{\sigma}_{eg} - \hat{a}\hat{\sigma}_{eg}^{\dagger}]
+ \sum_{m=L,R} i\hbar c\sqrt{\kappa/2\pi} \int dk \big[ \hat{b}_m(k)\hat{a}^{\dagger} - \hat{b}_m^{\dagger}(k)\hat{a} \big],
\]

where the first three terms are the free energy and the last term of the second line is the interaction between the atom and the cavity. The last term is the interaction between the cavity and the external fields, which are represented by operators \(\hat{b}_L\) and \(\hat{b}_R\) for the field modes on the left and right side of the cavity and \(\kappa\) is the total decay rate of the cavity field into the continuum. We assume a linear dispersion relation for the waveguides modes, where \(c\) is the velocity of the photons and \(\kappa\) is assumed to be frequency independent. The decay rate of the \(m\)th
level $\gamma_m$ represents both relaxation ($T_1$) processes and pure ($T_2$) dephasing (see supplemental material for details on the decoherence model [32]).

The protocol consists of two steps: first the interaction of the system with the gate photon, and second, sending in a multi-photon signal field. A key ingredient in our scheme is the difference in the scattering dynamics for photons scattering off the cavity conditioned on the state of the atom. If the atom is in state $|g\rangle$ it does not interact with the cavity field and the incoming resonant photon is transmitted through the cavity. On the other hand, if the atom is in state $|g\rangle$ the atom and the cavity form dressed states at frequencies $\omega_e \pm g$, where $\omega_e$ is the cavity frequency. An incoming photon at the undressed cavity frequency will be off-resonant and will be reflected. The state of the atom thus controls whether the cavity is transmitting the signal or not and thereby acts as a switch.

We can realize the single photon transistor if we can flip the state of the atom conditioned on the presence of a photon in the first gate pulse. To do this we use the procedure proposed in Ref. [23] and realized in Ref. [24]: We start with the atom initialized in the state $|\Psi_0\rangle = 1/\sqrt{2}(|g\rangle + |f\rangle)$ using the classical field on the transition between $|g\rangle$ and $|f\rangle$. Then we send in the gate pulse containing a single photon or not. By taking this gate pulse to be sent in symmetrically as shown in Fig. 2 (a) we ensure that we cannot determine the atomic state from whether the pulse has been transmitted or reflected. There is, however, a phase difference between the reflection and transmission such that in a suitable rotating frame the whole evolution is [23]

$$
|0\rangle |\Psi_0\rangle \rightarrow \frac{1}{\sqrt{2}} (|0\rangle (|g\rangle + |f\rangle))
$$

$$
|1\rangle |\Psi_0\rangle \rightarrow \frac{1}{\sqrt{2}} (|1\rangle (|g\rangle - |f\rangle)),
$$

where $|0\rangle$ ($|1\rangle$) denotes a zero (one) photon state of the field. Since the phase of the superposition now depends on the absence or presence of a single photon we can use a $\pi/2$-pulse on the $|f\rangle \rightarrow |g\rangle$ transition to map this onto the atomic population. The overall transformation in the first step is then given by

$$
|0\rangle |\Psi_0\rangle \rightarrow |0\rangle |g\rangle
$$

$$
|1\rangle |\Psi_0\rangle \rightarrow |1\rangle |f\rangle,
$$

which exactly realizes a conditional flip of the atomic state. The second step consists in sending in signal photons from the left as shown in Fig. 2 (b). Since the transmission or reflection of these photons is determined by the atomic state, the presence or absence of a photon in the first step controls the transmission in the second step. Hence the whole procedure implements a single photon transistor where the gate field controls the propagation of the large signal field for a time only limited by the lifetime $T_1$ of the atomic populations.

In the following we make a detailed analysis of the protocol in the presence of imperfections (extending the analysis of Ref. [23] where the protocol was mainly studied numerically). For the first step of the protocol we have a single incoming photon and we can expand the entire wave function on suitable basis states. The Schrödinger equation with the Hamiltonian [1] then results in two sets of differential equations for the amplitudes corresponding to the system starting in each of the uncoupled states $|g\rangle$ and $|f\rangle$. From solving the equations we find the reflection coefficient $r(k)$ for a photon with

![Fig. 1: Setup. (a) Three level ladder (Ξ) system coupled to a cavity field. The cavity field resonantly couples the upper states $|g\rangle$ and $|e\rangle$ of a Ξ-system with a coupling constant $g$ (single line), while it is far off resonance by a detuning $\Delta$ from the lower transition between $|f\rangle$ and $|g\rangle$, where it interacts with a coupling constant $g'$. The lower transition is driven by a classical field (double line). (b) Numerical and analytical results for the error probability for the first step of the protocol as a function of deviation from $\kappa = g$, parameterized by $\alpha = (\kappa - g)/g$. We use the parameters of 15 (coherence time $T_2^* = 2/(\gamma_g + \gamma_f) = 92$ $\mu$s, and anharmonicity $\Delta = (2\pi)206$ MHz), but a reduced coupling constant $g = (2\pi)5.3$ MHz. Full line: numerical results. Dashed line: analytical results. Dashed dotted line: approximate optimum discussed in the text. The pulse length is chosen according to an approximate optimization [32] with $n = 8$.](image1)

![Fig. 2: Steps of the single-photon transistor operation. (a) First step. A gate field is sent in symmetrically from the left and right. This makes it indistinguishable whether the photon was reflected or transmitted, but still a phase difference is imparted on the atom (see text). This phase difference can be used to flip the atom depending on the presence or absence of a photon in the field. (b) Second step. A strong signal field is sent in from the left. The field is reflected if the atom is in state $|g\rangle$ (full arrow), and transmitted if the atom is in state $|f\rangle$ (dashed arrow).](image2)
any pulse of temporal width $1$ the error probability only involves a higher order term the same group delay ensuring a maximal overlap, and $\kappa$ is the induced decay rate and coupling constant on the lower transition $[32]$. Experimentally the best coherence properties are currently reached for transmon qubits which have a rather low anharmonicity $\Delta \sim g$ $[18][20]$. For typical experimental parameters $g' \sim g$ and the induced decay would be highly detrimental. It will thus be desirable to reduce the coupling constant (e.g., by placing the atom near a field antinode) to suppress the off-resonant coupling. For a given superconducting system characterized by the anharmonicity $\Delta$ and the decoherence rates $\gamma_i$ we can then optimize the performance of the transistor by simultaneously optimizing the coupling constant and the pulse duration. Assuming $\kappa = g$ [including higher order terms omitted in Eq. (7)] we find that the optimal error scales as $P_{err} \sim (|\gamma_g + \gamma_f|/\Delta)^4/9$. In practice the condition $g - \kappa \approx 0$ may not be perfectly fulfilled due to fabrication imperfections etc. To characterize this we introduce a small parameter $\alpha = (\kappa - g)/g$, optimize the error probability assuming that we are limited by $\alpha$, and find $P_{err} \sim \alpha^2/3(|\gamma_g + \gamma_f|/\Delta)^1/9$. The full error will then be $\approx \max(P_1, P_2)$ (for further details see supplemental material $[32]$). These expressions show that the error rate can be quite low for realistic systems for which $\Delta \gg \gamma_i$. In particular, recent experiments with 3D confined transmon qubits have shown long coherence times. Taking parameters from Ref. $[18]$ we find that the error rate can be less than 1% for an imbalance $|\alpha| \lesssim 10\%$ (error 1.2% for the parameters of Ref. $[19]$). In Fig. 1 (b) we show the results of a direct numerical simulation of the first step of the protocol along with a detailed analytical theory and the approximate optimization. The results are seen to be in very good agreement. Alternatively, the proposal could also be implemented for transmon systems not confined to 3D cavities. Taking the parameters of Ref. $[20]$ we find an optimal error of $3.3\%$.

We now turn to the second step. Here we consider sending in a coherent state of constant amplitude for a time $T'$. This pulse is resonant with the cavity, but the transmission depends on the atomic state induced by the first signal pulse. The strongest difference in the distribution of the outgoing photons is obtained by sending in a strong beam resonant with the cavity with an intensity $I \gtrsim g^2/\kappa$. If the atom is in state $|f\rangle$ it essentially does not participate in the scattering dynamics and the cavity is just transmitting. If the atom is in state $|g\rangle$ the interaction shifts the cavity out of resonance and the light is reflected. Dealing with this situation in full detail is quite complicated and we therefore limit our calculations to a steady state calculation, which ignores dynamics on a time scale of $1/\kappa \ll T'$ (similar to the calculation in Ref. $[51]$). We numerically solve the density matrix equation in steady state for various incoming intensities, and

$$r(k) = \frac{\delta(k)(\delta(k) + \Gamma) - g^2}{(\delta(k) + \Gamma)(\delta(k) + i\kappa) - g^2}. \quad (4)$$

Here $\delta(k) = ck - \omega$ and $\Gamma = (\gamma_e - \gamma_g)/2$. This reflection coefficient is approximately unity in the regime we are interested in $g \gg \sqrt{\Gamma}, \sqrt{\kappa}, \delta$. Similarly, the transmission amplitude is found to be

$$t(k) = \frac{-i\kappa(\delta(k) + i\Gamma)}{(\delta(k) + i\Gamma)(\delta(k) + i\kappa) - g^2}, \quad (5)$$

which vanishes in the same limit. When the atom is in state $|f\rangle$ the transmission and reflection coefficient can be found by setting $g = \Gamma = 0$ in the above expression and for $\kappa \gg \delta$ these can be then approximated to be $r = 0$ and $t = 1$. These results justify the evolution in $[2]$ provided that the bandwidth of the pulse is smaller than $\min(\kappa, g^2/\kappa)$. [Above we have rescaled all amplitudes by $\exp(\gamma_g t/2)$ and $\exp(\gamma_f t/2)$ to remove the damping of the levels $|g\rangle$ and $|f\rangle$].

To get a concrete estimate of the functioning of the transistor we assume the incident gate pulse to have a Gaussian temporal profile

$$f(t) \propto e^{-(t - T/2)^2/4\sigma_T^2}, \quad (6)$$

with a width $\sigma_T$. We restrict the dynamics to a finite time interval $[0, T]$ and choose the ratio $n = T/\sigma_T$ sufficiently big, e.g., $n = 8$, that we can ignore the pulse area outside this interval. Assuming equal probabilities to receive a photon or not during the first step we find the average error probability for setting the qubit cavity system

$$P_{err} \approx \frac{1}{4} \left( (\gamma_g + \gamma_f) n\sigma_T + \frac{\kappa(\gamma_e - \gamma_g)}{g^2} + \frac{1}{2\sigma_T^2\kappa^2} \left( 1 - \frac{\kappa^2}{g^2} \right)^2 \right). \quad (7)$$

This result is valid in the limit where all terms are small. It accounts for three sources of error: The first term is given by the decoherence of the levels $|g\rangle$ and $|f\rangle$. The second term represents the fact that the cavity is not ideal (since this term accounts for decoherence while in state $|e\rangle$ and the full decoherence of level $|g\rangle$ is included in the first term, it is only the additional decoherence $\gamma_e - \gamma_g$ which appears here). The third term is the error due to the finite bandwidth of the pulse which causes an imperfect phase shift with the atom. This term is minimized for $\kappa = g$, which ensures that the dispersion relation for photons at the central frequency is the same regardless of whether the atom is in state $|f\rangle$ or $|g\rangle$. At $\kappa = g$ the outgoing photon wave packets will thus have the same group delay ensuring a maximal overlap, and the error probability only involves a higher order term $1/(\sigma_T)^6$. With this choice the transistor is efficient for any pulse of temporal width $1/g \lessgtr \sigma_T \lesssim 1/n(\gamma_g + \gamma_f)$.

So far we have assumed that the coupling between the two lower states can be neglected. The detuning $\Delta$ of the cavity from the lower transition equals the anharmonicity of the atom and for large $\Delta$ the coupling can be taken into account by adding an additional error to the first step $P_{err} \approx \ldots + 4\kappa n \sigma_T$, where $\kappa = g^2/\Delta^2 + \kappa^2$ and $g'$ is the induced decay rate and coupling constant on the lower transition $[32]$. Experimentally the best coherence properties are currently reached for transmon qubits which have a rather low anharmonicity $\Delta \sim g$ $[18][20]$. For typical experimental parameters $g' \sim g$ and the induced decay would be highly detrimental. It will thus be desirable to reduce the coupling constant (e.g., by placing the atom near a field antinode) to suppress the off-resonant coupling. For a given superconducting system characterized by the anharmonicity $\Delta$ and the decoherence rates $\gamma_i$ we can then optimize the performance of the transistor by simultaneously optimizing the coupling constant and the pulse duration. Assuming $\kappa = g$ [including higher order terms omitted in Eq. (7)] we find that the optimal error scales as $P_{err} \sim (|\gamma_g + \gamma_f|/\Delta)^4/9$. In practice the condition $g - \kappa \approx 0$ may not be perfectly fulfilled due to fabrication imperfections etc. To characterize this we introduce a small parameter $\alpha = (\kappa - g)/g$, optimize the error probability assuming that we are limited by $\alpha$, and find $P_{err} \sim \alpha^2/3(|\gamma_g + \gamma_f|/\Delta)^1/9$. The full error will then be $\approx \max(P_1, P_2)$ (for further details see supplemental material $[32]$). These expressions show that the error rate can be quite low for realistic systems for which $\Delta \gg \gamma_i$. In particular, recent experiments with 3D confined transmon qubits have shown long coherence times. Taking parameters from Ref. $[18]$ we find that the error rate can be less than 1% for an imbalance $|\alpha| \lesssim 10\%$ (error 1.2% for the parameters of Ref. $[19]$). In Fig. 1 (b) we show the results of a direct numerical simulation of the first step of the protocol along with a detailed analytical theory and the approximate optimization. The results are seen to be in very good agreement. Alternatively, the proposal could also be implemented for transmon systems not confined to 3D cavities. Taking the parameters of Ref. $[20]$ we find an optimal error of $3.3\%$.

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evaluate the reflected and transmitted intensities from the steady state solution [32]. With this we find that for \( \kappa \sim g \) the maximal reflected intensity \( I_R \) saturates in the range \( I_R = g^2/4\kappa \) to \( I_R = g^2/2\kappa \) for strong incoming intensities \( I_{IN} \gg g^2/\kappa \) and small atomic decay \( \gamma_i \ll \kappa \) as show in Fig. 3. The same conclusion is reached from a perturbative calculation valid for strong driving (see supplemental material [32]). Taking the lower of these two values, the gain of the transistor \( G \), defined as the difference in the reflected number of photons condition on an incident photon, is given by \( G = T'g^2/4\kappa \). The time \( T' \) is limited by the decay of state \( |g\rangle \) to \( |f\rangle \). The largest average gain is found for \( T' \geq T_1 \) where we obtain \( G = T_1g^2/4\kappa \). Taking the optimized settings derived for the first step of the protocol we find gains of \( G \approx 783, \) 570, and 116 for the numbers of Refs. [18], [19], and [20], respectively. This last step represents the slowest part of the procedure which thus limits the duty cycle of the transistor, but the speed can be increased by reducing \( T' \) at the expense of also reducing the gain.

![FIG. 3: (a) Reflected intensity \( I_R \) for an atom in state \( |g\rangle \) as a function of incoming intensity \( I_{IN} \) obtained from a numerical solution of the master equation for various values of \( g/\kappa \). For weak incoming intensities the reflected intensity is equal to the total incoming intensity, but saturates in the range \( g^2/4\kappa \) (dotted line) to \( g^2/2\kappa \) as the intensity is increased. We have assumed \( \kappa = g \), and that we only have population decay of the upper level at a rate \( \gamma_i = 1.3 \cdot 10^{-7} \) corresponding to the values in Fig. 1 with \( \gamma_i = 2g \). The saturation of the reflected intensity provides an upper limit for how well one can distinguish an atom being in \( |g\rangle \) from an atom being in \( |f\rangle \) which has a vanishing reflection. This saturation thus limits the gain of the transistor. (b) Alternative implementation based on a two-level system coupled to a cavity. Taking the subsystem to have a vacuum Rabi coupling \( g \), the nonlinearity of the Jaynes-Cummings Hamiltonian can mimic our example three-level system. Denoting the atomic levels by \( \tilde{g}, \tilde{e} \) and the subsystem cavity Fock states with \( \tilde{0}, \) etc., we identify \( |f\rangle = \langle \tilde{g}, 0 | - \langle \tilde{e}, 0 | \rangle/\sqrt{2}, |e\rangle = \langle \tilde{g}, 2 | - \langle \tilde{e}, 2 | \rangle/\sqrt{2}. \) This yields an anharmonicity parameter \( \Delta = 2g(1 - \sqrt{2}) \), while the coupling for the effective three-level system \( g \) is determined by the coupling of the anharmonic cavity-subsystem to another cavity and can be as large as desired by appropriate cavity design.

In summary, we have shown that using superconducting circuit QED it is possible to realize a working photonic transistor sensitive to only a single incoming microwave photon. The device can readily be implemented with existing technology. We have focussed on an implementation based on a three level artificial atom but alternatively the procedure could also be implemented in a two-level system in the strong dispersive regime [19, 25] or by using a setup similar to [26] to mimic the three-level atom by a two-level atom strongly coupled to an optical cavity as outlined in Fig. 3 (b), opening up our transistor technique to a wide variety of additional systems. An immediate application of the device proposed here would be the efficient quantum non-demolition detection of individual itinerant microwave photons. In addition, the device could easily be incorporated into many quantum information protocols working as an efficient processor of information encoded in the photon field. As a particular example, for higher incident photon numbers the device will serve as a parity detector, detecting the parity of the incident photon number [27]. This is an important primitive for several suggested continuous variable quantum information protocols [28, 29].

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[1] D. E. Chang, A. S. Sørensen, A.Demler, and M. D. Lukin, Nat. Phys. 3, 807 (2007).
[2] W. Chen, K. M. Beck, R. Bücker, M. Gullans, M. D. Lukin, H. Tanji-Suzuki, and V. Vuletić, Science 341, 768 (2013).
[3] A. V. Akimov, A. Mukherjee, C. L. Yu, D. E. Chang, A. S. Zibrov, P. R. Hemmer, H. Park, and M. D. Lukin, Nature 450, 402 (2007).
[4] J. Claudon, J. Bleuse, N. S. Malik, M. Bazin, P. Jaffrennou, N. Gregersen, C. Sauvan, P. Lalanne, and J. M. Gerard, Nat. Phot. 4, 174 (2010).
[5] K. G. Lee, X. W. Chen, H. Eghlidi, P. Kukura, R. Lettow, A. Renn, V. Sandoghdar, and S. Götzinger, Nat. Phot. 5, 166 (2011).
[6] T. Wilk, S. C. Webster, A. Kuhn, and G. Rempe, Science 317, 488 (2007).
[7] T. Volz, A. Reinhard, M. Winger, A. Badolato, K. J. Hennessy, E. L. Hu, and A. Imamoğlu, Nat. Phot. 6, 605 (2012).
[8] D. Englund, A. Majumdar, M. Bajcsy, A. Faraon, P. Petroff, and J. Vuckovic, Phys. Rev. Lett. 108, 093604 (2012).
[9] R. Maiwald, D. Leibfried, J. Britton, J. C. Bergquist, G. Leuchs, and D. J. Wineland, Nat. Phys. 5, 551 (2009).
[10] R. J. Schoelkopf and S. M. Girvin, Nature 451, 664 (2008).
[11] A. Wallraff, D. I. Schuster, A. Blais, L. Frunzio, R.-S. Huang, J. Majer, S. Kumar, S. M. Girvin, R. J. and Schoelkopf, Nature 431, 162 (2004).
[12] S. Rebic, J. Twamley, and G. J. Milburn, Phys. Rev. Lett. 103, 150503 (2009).
[13] L. Neumeier, M. Leib, and M. J. Hartmann, Phys. Rev. Lett. 111, 63601 (2013).
[14] P. Adhikari, M. Hafezi, and J. M. Taylor, Phys. Rev. Lett. 110, 060503 (2013).
[15] D. Ristè, M. Dukalski, C. A. Watson, G. de Lange, M. J. Tiggelman, Ya. M. Blanter, K. W. Lehnert, R. N. Schouten, and L. DiCarlo, cond-mat/1306.4002.
[16] J. Koch, T. M. Yu, J. Gambetta, A. A. Houck, D. I. Schuster, J. Majer, A. Blais, M. H. Devoret, S. M. Girvin, and R. J. Schoelkopf, Phys. Rev. A 76, 042319 (2007).
[17] A. Fedorov, L. Steffen, M. Baur, M. P. da Silva, and A. Wallraff, Nature 481, 170 (2012).
[18] C. Rigetti, J. M. Gambetta, S. Poletto, B. L. T. Plourde, J. M. Chow, A. D. Córcoles, J. A. Smolin, S. T. Merkel, J. R. Rozen, G. A. Keefe, M. B. Rothwell, M. B. Ketchen, and M. Steffen, Superconducting qubit in waveguide cavity with coherence time approaching 0.1 ms, Phys. Rev. B 86, 100506 (2012).
[19] A. P. Sears, A. Petrenko, G. Catelani, L. Sun, H. Paik, G. Kirchmair, L. Frunzio, L. I. Glazman, S. M. Girvin, and R. J. Schoelkopf, Phys. Rev. B 86, 180504 (2012).
[20] A. A. Houck, J. Koch, M. H. Devoret, S. M. Girvin, R. J. Schoelkopf, Quan. Inf. Proc. 8, 105 (2009).
[21] M. Neeley, R. C. Bialczak, M. Lenander, E. Lucero, M. Mariantoni, A. D. O’Connell, D. Sank, H. Wang, M. Weides, J. Wenner, Y. Yin, T. Yamamoto, A. N. Cleland, and J. M. Martinis, Nature 467, 570 (2010).
[22] D. Witthaut and A. S. Sørensen, New J. Phys. 12, 043052 (2010).
[23] L.-M. Duan and H. J. Kimble, Phys. Rev. Lett. 92, 127902 (2004).
[24] H. Kim, R. Bose, T. C. Shen, G. S. Solomon, and E. Waks, A quantum logic gate between a solid-state quantum bit and a photon, Nat. Phot. 7, 373 (2013).
[25] M. Hatridge, S. Shankar, M. Mirrahimi, F. Schackert, K. Geerlings, T. Brecht, K. M. Sliwa, B. Abdo, L. Frunzio, S. M. Girvin, R. J. Schoelkopf, and M. H. Devoret, Science 339, 178 (2013).
[26] K. Stannigel, P. Komar, S. J. M. Habraken, S. D. Bennett, M. D. Lukin, P. Zoller, and P. Rabl, Phys. Rev. Lett. 109, 013603 (2012).
[27] S. R. Sathyamoorthy, L. Tornberg, A. F. Kockum, B. Q. Baragiola, J. Combes, C. M. Wilson, T. M. Stace, and G. Johansson, quant-ph/1308.2208.
[28] S. L. Braunstein and P. van Loock, Rev. Mod. Phys. 77, 513 (2005).
[29] K. Banaszek and K. Wódkiewicz, Phys. Rev. A 58, 4345 (1998).
[30] P. Meystre, M. Sargent III, Elements of Quantum Optics, (Springer, New York, 1999).
[31] C. J. Hood, M. S. Chapman, T. W. Lynn, and H. J. Kimble, Phys. Rev. Lett. 80, 4157 (1998).
[32] See supplemental material.