Generation of Nonclassical Microwave States Using an Artificial Atom in 1D Open Space

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We have embedded an artificial atom, a superconducting transmon qubit, in a 1D open space and investigated the scattering properties of an incident microwave coherent state. By studying the statistics of the reflected and transmitted fields, we demonstrate that the scattered states can be nonclassical. In particular, by measuring the second-order correlation function, $g^{(2)}$, we show photon antibunching in the reflected field and superbunching in the transmitted field. We also compare the elastically and inelastically scattered fields using both phase-sensitive and phase-insensitive measurements.

A single atom interacting with propagating electromagnetic fields in open space is a fundamental system of quantum optics. Strong coupling between a single artificial atom and resonant propagating fields has recently been achieved in a 1D system [1,2], experimentally demonstrating nearly perfect extinction of the forward propagating fields [2]. However, this extinction can be explained by classical theory: a classical pointlike oscillating dipole perfectly reflects resonant incident fields [3]. In this Letter, we demonstrate the quantum nature of the scattered field generated from our artificial atom in 1D open space by using a resonant coherent state as the incident field. In particular, by measuring the second-order correlation function, we show the reflected field is antibunched [4,5] while still maintaining first-order coherence. Moreover, we observe superbunching statistics in the transmitted fields [4].

To understand how our artificial atom generates antibunched and superbunched states, it is helpful to consider the incident coherent state in the photon number basis. For a low power incident field with less than 0.5 average photons per lifetime of our atom, we can safely approximate the coherent field using only the first three photon eigenstates. If we consider a one-photon incident state, the atom reflects it, leading to antibunching statistics in the reflected field. Together with the zero-photon state the reflected field still maintains first-order coherence. For a two-photon incident state, since the atom is not able to scatter more than one photon at a time, the pair has a much higher probability of transmission, leading to superbunching statistics in transmission [4,6]. In this sense, our single artificial atom acts as a photon-number filter, which extracts the one-photon number state from a coherent state. This represents a novel way to generate photon correlations and nonclassical states at microwave frequencies compared with other recent work [7–11].

Our system consists of a superconducting transmon qubit [12], strongly coupled to a 1D coplanar waveguide transmission line [see Fig. 1(a)]. The ground state $|0\rangle$ and first excited state $|1\rangle$ have a transition energy $\hbar\omega_{01}$. The relaxation rate of the qubit is dominated by an intentionally strong coupling to the 50 Ω transmission line through the coupling capacitor $C_c$, as shown in Fig. 1(b).

The electromagnetic field in the transmission line is described by an incoming voltage wave $V_{in}$, a reflected wave $V_R$, and a transmitted wave $V_T$. In Fig. 1(a), the transmittance is defined as $T = |V_T/V_{in}|^2$. For a weak coherent drive on resonance with the atom, we expect to see full reflection of the incident signal [4,13]. This can be understood in terms of interference between the incident wave and the wave scattered from the atom, which destructively interfere in transmission and constructively interfere in reflection [4,13]. In the sample measured here, we achieved extinction of more than 99% in transmittance, as shown in Fig. 1(c). By measuring the transmission coefficient as a function of probe frequency and probe power $P$, we extract $\omega_{01}/2\pi = 5.12$ GHz, $\Gamma_{10}/2\pi = 41$ MHz, and $\Gamma_\phi/2\pi = 1$ MHz [2]. The relaxation rate $\Gamma_{10}$ is dominated by coupling to the transmission line and

![FIG. 1 (color online).](image)

(A) A micrograph of our artificial atom, a superconducting transmon qubit embedded in a 1D open transmission line. (Zoom In) Scanning-electron micrograph of the SQUID loop of the transmon. (B) Schematic setup for measurement of the second-order correlation function. This setup enables us to do Hanbury–Brown–Twiss measurements between output ports 1 and 2. Depending on the choice of input port, we can measure $g^{(2)}$ of the reflected or transmitted field. (C) Transmittance on resonance as a function of incident power. (Inset) A weak, resonant coherent state is reflected by the atom.
is much greater than the pure dephasing rate $\Gamma_\phi$ in our system. We define an average number of photons per interaction time $2\pi/\Gamma_{10}$ as $N = 2\pi P/(\hbar \omega_p \Gamma_{10})$. $N = 1$ for a power of $-128$ dBm.

The resonant electromagnetic field reflected from or transmitted through the atom (depending on the choice of input port, see Fig 1(b)) is fed through two circulators to a commercial 90° hybrid coupler, with its other input terminated with 50 $\Omega$. The hybrid coupler effectively acts as a microwave beam splitter. Ideally, the signal coming into the other input of the hybrid coupler should be vacuum. The two outputs of the beam splitter are sent to two nominally identical microwave HEMT amplifiers at 4.2 K that have system noise temperatures of 7 K. We make the assumption in our analysis that the noise added by the two amplifiers is uncorrelated. After further amplification, the two outputs are fed into a pair of vector digitizers, which capture the voltage amplitude. We then can choose to digitally filter the voltage data to a desired bandwidth (BW). This setup enables us to perform Hanbury–Brown–Twiss (HBT) [14–18] measurements using linear digitally filter (average) the data in all the $g^{(2)}$ measurements. For a coherent state, we expect $g^{(2)}(t) = 1$. This is indeed what we find if our atom is off-resonance from our applied coherent source.

After these initial measurements, we measured second-order correlations of the field transmitted through our qubit. We applied an on-resonance microwave drive and measured $g^{(2)}(\tau)$ in transmission for different incident powers as shown in Fig. 2(b). Here, the sampling frequency was again set to $10^8$ samples/s with BW = 55 MHz. At the lowest power ($P = -129$ dBm, $N = 0.8$) that we can readily measure, we see superbunching of the photons [4], with $g^{(2)}(\tau = 0) = 2.31 \pm 0.09 > 2$. We see superbunching even with our measured $g^{(2)}(0)$ suppressed as a result of trigger jitter [see Fig. 2(e)]. Superbunching occurs because the one-photon state of the incident field has been selectively filtered and reflected while the two-photon state is more likely transmitted (the three-photon components and higher are negligible). This transmitted state generated from our qubit is thus bunched even more than a thermal state. For high powers, where $N \gg 1$, we find $g^{(2)}(\tau) = 1$, as most of the coherent signal passes through the transmission line without interacting with the atom owing to saturation of the atomic response. The correlation measurements once again resemble those of a coherent state. For all measurements shown here we find, $g^{(2)}(\infty) = 1$, as expected. In the inset of Fig. 2(b), we plot $g^{(2)}(0)$ as a function of incident power and clearly see the bunching behavior decrease as the incident power increases. For comparison, we also plot $g^{(2)}(0)$ for a coherent state and thermal state.

In Fig. 2(d), we plot the measured $g^{(2)}(\tau)$ of the reflected field from our atom. At low powers, where $N \ll 1$, we clearly observe antibunching of the field [4]. Each trace here was collected, computed, and averaged over 17 hours, corresponding to $2.4 \times 10^{11}$ measured quadrature field samples (2 Tbyte of data). The antibunching behavior at the lowest power ($P = -132$ dBm, $N = 0.4$), $g^{(2)}(0) = 0.51 \pm 0.05$, reveals the quantum nature of the field [5]. Ideally, we would find $g^{(2)}(0) = 0$ as the atom can only absorb and emit one photon at a time. The nonzero $g^{(2)}(0)$ we measured originates from four effects: (1) a thermal field at 50 mK, (2) a finite filter bandwidth (here 55 MHz), (3) trigger jitter between the two digitizers, and (4) stray fields including background reflections in the line and leakage through circulator 1 [Fig. 1(b)]. In Fig. 2(e), we measure $g^{(2)}(\tau)$ at $P = -131$ dBm for different filter bandwidths and clearly see that the antibunching dip...
depends on BW. For a small BW, i.e., long sampling time, the time dynamics of antibunching cannot be resolved. In other words, within the sampling time, the atom is able to absorb and emit multiple photons. If BW \( \ll \Gamma_{10} \), \( \Omega_p \), where \( \Omega_p \) is the Rabi frequency, the antibunching dip we measure vanishes entirely. This interplay between BW and \( \Omega_p \) yields a power dependent \( g^{(2)}(0) \), as shown in the inset of Fig. 2(d).

In Fig. 2(f), we show how all four factors listed above combine to produce the theoretical curves for our measured \( g^{(2)}(\tau) \). The partial theory curves include finite temperature and filter bandwidth, but not leakage and jitter. The green (no leakage) curve includes everything but (4) and the red curve (complete theory) includes all four effects. The solid lines in Fig. 2(b)–2(f) are theoretical results based on a master equation formalism. The digital filter is modeled by a single-mode resonator. A master equation describing both the transmon and the resonator is derived using the formalism of cascaded quantum systems. To model the effect of the trigger jitter, the value of \( g^{(2)}(\tau) \) at each point is replaced by the average value of \( g^{(2)}(\tau), g^{(2)}(\tau - 10 \text{ ns}) \) and \( g^{(2)}(\tau + 10 \text{ ns}) \).

Our artificial atom selectively filters out the Fock state \( n = 1 \) from the input coherent state. As a result, the reflected and transmitted field display antibunched and superbunched statistics, respectively. Thus, the qubit acts as a passive photon-number filter, converting a coherent microwave state to a nonclassical one, with high production rate.

While the scattered field requires a purely quantum description, it can still maintain first-order coherence similar to a classical field, as shown below. We can define the first-order correlation function in steady state as \( g^{(1)} = \langle V^2 \rangle / \langle V \rangle^2 \). First-order coherence then refers to \( g^{(1)} = 1 \). For a thermal source this function is 0 and for a coherent state it is 1.
The first-order coherence properties of the scattered resonance field strongly depend on the Rabi frequency $\Omega_p$ and the relaxation rate $\Gamma_{10}$ of the atom. The Rabi frequency $\Omega_p$ is linearly proportional to the amplitude of the drive, $\Omega_p \propto \sqrt{P}$ [2]. For $\Omega_p \ll \Gamma_{10}$, we expect the scattered fields to be coherent to first order, with a power spectrum well described by a delta function at $\omega_{01}$ [5]. For $\Omega_p \gg \Gamma_{10}$, the scattered fields contain three additional inelastic Lorentzians (known as the Mollow Triplet [20]).

We send a single tone at $\omega_{01}$ and measure these scattered (reflected) fields from only one of the output ports, as shown in Fig. 3(a). Note that we see the same behavior from both outputs. We use a phase-sensitive average $\langle V^2 \rangle$ to capture the elastic (coherent) component of the scattered field. For the total scattered field, the sum of the elastically and inelastically scattered fields, we use a phase-insensitive average $\langle V^2 \rangle$. The amount of the inelastic field that we capture depends on the bandwidth of our measurement, as indicated in Fig. 3(a). The solid curves are the theory fits, using the model in Fig. 3(a) (integrating the Mollow triplet), with no free-fitting parameters (using the same parameters extracted before). As expected, at low incident power, the total scattered field is roughly equal to the elastically scattered (coherent) field. This indicates that the field is first-order coherent with $g^{(1)} \approx 1$. We note that this is also the regime where antibunching is observed. At high incident fields, where $\Omega_p > \Gamma_{10}$, the main contribution to the total field is from inelastic scattering.

In conclusion, we investigated the scattering properties of a single artificial atom in a 1D open space by measuring the first, second, and fourth moments of the voltage field. We verified the quantum nature of the scattered field using the second-order correlation function, while also showing that the field maintained first-order coherence. In fact, the whole process leads to a redistribution of the photon number state [6]. We plan to investigate applications of this phenomenon, such as generating single-photon states on demand. This system may offer advantages over placing an artificial atom in a cavity [21–23]. For instance, the generated single photons can maintain the same envelope as the incident coherent state with a wide bandwidth limited only by the atom relaxation rate. For a cavity-based photon source, the bandwidth is limited by the cavity width and subject to the problem of stochastic release by the cavity.

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[19] Technically, it is a chaotic state.


