Vacuum-induced Autler-Townes splitting in a superconducting artificial atom

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We experimentally study a vacuum-induced Autler-Townes doublet in a superconducting three-level artificial atom strongly coupled to a coplanar waveguide resonator and simultaneously to a transmission line. The Autler-Townes splitting is observed in the reflection spectrum from the three-level atom in a transition between the ground state and the second excited state when the transition between the two excited states is resonant with a resonator. By applying a driving field to the resonator, we observe a change in the regime of the Autler-Townes splitting from quantum (vacuum-induced) to classical (with many resonator photons). Furthermore, we show that the reflection of propagating microwaves in a transmission line could be controlled by different frequency single photons in a resonator.

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Electromagnetic waves propagating through a medium of identical atoms are resonantly absorbed. The absorption can be eliminated when a strong driving field couples other atomic transitions, for example, three-level atoms, creating a transparency window for the waves. This results in either Autler-Townes splitting (ATS) or electromagnetically induced transparency (EIT) and is extensively studied in quantum optics. The same phenomena can be observed even if the medium is replaced by a single atom or a molecule which is coupled to either a cavity or an open space. However, the strong coupling between the driving field and a single natural atom is difficult to achieve. Recently, the strong coupling has been experimentally achieved between a superconducting artificial atom and non-quantized microwave fields confined in a one-dimensional transmission line. This enables ATS and EIT to be observed using a single superconducting artificial atom. Several experiments have already shown ATS and state manipulation in a single three-level artificial atom. ATS and EIT have been demonstrated in a Jaynes-Cummings ladder system of a single two-level artificial atom dispersively coupled to a cavity.

The observation of both ATS and EIT in a medium (an ensemble of atoms) usually requires strong classical driving fields. However, the conditions are different: the driving field required for the observation of EIT is weaker than that for ATS, and there is a crossover from ATS to EIT. In particular, in ATS the splitting is determined by the Rabi frequency of the driving field, whereas in EIT the dip between the two peaks is a result of quantum interference. These effects take place even if the medium is scaled down to a single atom and the driving field is quantized. The condition for the observation of the effects can then be reduced to a small number of photons or even without photons owing to coupling to a vacuum mode. Theoretical investigations show that the transparency for the classical probing field can still occur. This has potentially important applications, for example, in single-photon switches and transistors, all-optical quantum logic, quantum communication, and metrology. However, vacuum-induced transparency has only recently been observed in an ensemble of three-level atoms. Here, we demonstrate ATS with a single three-level artificial atom, which is controlled by a quantized single-mode field in a transmission line resonator.

Our device, shown in Fig. 1(a), consists of a three-level artificial atom capacitively and strongly coupled to two macroscopic objects simultaneously: a coplanar waveguide resonator (CPWR) and a transmission line. Although the device is reminiscent of the one studied in Ref. [22], where the atom was coupled to two transmission lines, it is essentially different. We emphasize that in spite of the strong coupling of the atom to the transmission line and the atom to the resonator, the line and resonator are decoupled from each other (details are given...
in Supplementary Material [33]). This is a peculiarity of the device owing to the unique property of superconducting quantum systems: they are micrometer-scale electronic circuits. Namely, the two macroscopic objects (i.e., open lines, resonators) can be effectively decoupled from each other but, at the same time, strongly coupled to the quantum circuit. The artificial atom has the geometry of a tunable-gap superconducting flux qubit: a superconducting loop with two Josephson junctions and a dc-SQUID (α-loop) [33,34], shown in Fig. 1(b), is fabricated near the voltage antinode (open end) of the CPWR using electron-beam lithography and Al/AIO$_x$/Al shadow evaporation (see the SEM image in [33]). A weak probe signal with frequency $\omega_p$ is applied from the transmission line (left-hand side). Both the half-wavelength CPWR and the transmission line are etched from a 50 nm thick Nb thin film sputtered on an oxidized undoped silicon wafer. The lines have a center conductor of 10 μm width separated from the ground planes by 6 μm gaps, resulting in a 50 Ω wave impedance.

Our experiment is carried out in a dilution refrigerator with a base temperature of about 30 mK. The fundamental frequency of the CPWR is $\omega_p = 2\pi \times 8.794$ GHz with the decay rate $\kappa = 2\pi \times 3.6$ MHz. The three lowest energy levels of the artificial atom, denoted by $|g\rangle$ for the ground state, $|e\rangle$ for the first excited state, and $|f\rangle$ for the second excited state, are controlled by the magnetic flux $\Phi$ in the loop. The minimum transition energies from the ground to both excited states occur at half-integer flux quanta $\Phi_N = (N + \frac{1}{2})\Phi_0$, where $N$ is an integer and $\Phi_0$ is the flux quantum. The minimal energy gap $\Delta$ for the transition from $|g\rangle$ to $|f\rangle$ can be tuned by controlling the magnetic flux $\Phi_n$ penetrating through the α-loop. When the biased flux $\Phi$ is close to $\Phi_N$, the two lowest-energy eigenstates ($|g\rangle$ and $|e\rangle$) are in superposition of the two oppositely circulating persistent current states of the loop. The energy between the two lowest levels $h\omega_{ge}$ in the vicinity of $\Phi_N$ is approximated by the expression $\sqrt{(21p\delta\Phi)^2 + \Delta(\Phi_N)^2}$, where $\delta\Phi = \Phi - \Phi_N$ and $I_p$ is the persistent current in the qubit loop. The weak dependence of $\Delta$ on $\Phi$ can be safely neglected when $\delta\Phi \ll \Phi_0$, which means that $\Delta(\Phi) \approx \Delta(\Phi_N)$. As shown in Fig. 1(c), by choosing $\Phi_N$ and $\delta\Phi$, we can tune both transition frequencies $\omega_{ef}$ and $\omega_{gf}$. The ATS can be observed by measuring directly the reflection spectrum through the left transmission line.

As shown in Fig. 1(c), the transition energy between the $|e\rangle$ and $|f\rangle$ states is aligned to the resonator resonance. The probe field is applied at the $|g\rangle$ to $|f\rangle$ transition. Thus, the effective Hamiltonian of the whole system driven by classical waves with amplitude $\Omega$ at frequency $\omega_p$ close to the $|g\rangle \leftrightarrow |f\rangle$ transition is given by

$$
H = h\omega_e a^\dagger a + h\omega_{gf} \sigma_{ff} + h\omega_{ge} \sigma_{ee} + h\Delta_0 \sigma_{ff} + 2h\Omega (\sigma_{ff} + \sigma_{gf}) \cos \omega_p t.
$$

(1)

Here, the atomic operator $\sigma_{jk} = |j\rangle\langle k|$ with $\{|j\rangle\} = \{|g\rangle, |e\rangle, |f\rangle\}$ and $a^\dagger$ and $a$ are the photon creation and annihilation operators in the single-mode resonator, respectively. The first four terms of Eq. (1) represent the Jaynes-Cummings Hamiltonian for the two-level system with the following specific features: it presents the interaction of two excited states ($|e\rangle$ and $|f\rangle$) with the resonator rather than the ground and excited states. Nevertheless, the Jaynes-Cummings physics can be applied here [35,37]. Namely, in the vicinity of the $|e\rangle \leftrightarrow |f\rangle$ resonance with the resonator, $|f\rangle$ level splits due to the interaction with the resonator vacuum mode, which can be described in terms of zero-photon dressed states as $|v_0\rangle = \cos \theta |f0\rangle - \sin \theta |e1\rangle$ and $|u_0\rangle = \sin \theta |f0\rangle + \cos \theta |e1\rangle$, where $\tan 2\theta = 2g_{0}/(\omega_r - \omega_f + \omega_e)$ and the state $|n\rangle = |j\rangle \otimes |n\rangle$ with the two quantum numbers $j$ and $n$, denoting the three-level atomic states $|j\rangle$ and the Fock states of resonator $|n\rangle = \{|0\rangle, |1\rangle, |2\rangle, \ldots \}$, respectively. The dressed states are schematically shown in the left panel of Fig. 1(d), while the right panel shows the general case of a system with $n$ photons in the resonator. The last term represents the probing field which couples the ground and the second excited states ($|g\rangle \leftrightarrow |f\rangle$), where the vacuum mode $|0\rangle$ is omitted for
The anticrossing measured in the transition $|e\rangle \leftrightarrow |f\rangle$ is visible due to the thermal population of the first excited state $|e\rangle$.

After confirming the vacuum-induced ATS, we study the dependence of the ATS on the resonant driving to the resonator field at $\omega_p \approx 2\pi \times 8.794$ GHz using a probe signal at $\omega_p \approx 2\pi \times 12.2$ GHz, which corresponds to the transition between the ground state and the second excited state $(|g\rangle \leftrightarrow |f\rangle)$. The dynamics of this driven atom-resonator system can be described by the Markovian master equation:

$$\dot{\rho} = -i\frac{\hbar}{\gamma} [H_F, \rho] + \kappa (2\sigma_a^+ \rho \sigma_a^- - \sigma_a^+ \rho \sigma_a^- + \sigma_a^- \rho \sigma_a^+) + \sum_{j,k=0,e,f} \frac{\gamma_{jk}}{2} [2\sigma_{jk} \rho \sigma_{kj} - \sigma_{jk} \sigma_{kj} \rho - \rho \sigma_{jk} \sigma_{kj}],$$

for the density matrix $\rho$ of the coupled system, where $\gamma_{jk}$ is the relaxation/dephasing rate of the atom. Here, $H_F = H + \hbar \Omega_d (a^+ e^{-i\omega_{at}} + a e^{i\omega_{at}})$ is the Hamiltonian when the driving with amplitude $\Omega_d$ to the cavity field is considered and the Hamiltonian $H$ is given by Eq. (1).

The reflection coefficient versus frequency is shown in Fig. 2(a), where the splitting amplitude to be $\Omega/2\pi \approx 7$ MHz.
Figure 3: (a) Dependence of ATS on driving power of the resonator. The orange guidelines show the splitting. The inset shows dip positions versus calibrated photon numbers. The solid lines are fits accounted the ac Zeeman shift. (b) Simulated ATS in the three-level atom coupled to a resonator. (c) Observed splittings at $\langle n \rangle = 0$ (quantum regime), $\langle n \rangle = 1$ and $\langle n \rangle = 5$ (upper three panels). The solid lines show the results of calculations. The bottom panel shows the difference between the traces with $\langle n \rangle = 1$ and $\langle n \rangle = 0$. (d) The splitting at $P = -113$ dBm, corresponding to $\langle n \rangle \approx 35$.

onator with the probing amplitude $\Omega/2\pi \approx 7$ MHz for the $|g\rangle \leftrightarrow |f\rangle$ transition close to the weak driving limit $[33]$. In this regime, the splitting is determined by the coupling to the vacuum mode of the resonator. The vacuum-induced ATS is a quantum mechanical phenomenon, which cannot be observed in the classically driven ATS. At larger power, the splitting starts to increase according to $2g_0\sqrt{\langle n \rangle} + 1$, where $\langle n \rangle$ is the mean photon number inside the resonator. We call quantum regime the one, where the vacuum splitting gives the main contribution to the ATS. Note that there are other quantum effects that can be observed beyond the subject of the present work (e.g., resolving photon states due to quantized fields in the resonator). The asymmetry of the branches is explained by the ac Zeeman shift (or ac Stark shift) $[32,33]$ from the $|g\rangle \leftrightarrow |f\rangle$ transition dispersively coupled to the resonator, and the frequency shift $\langle \zeta n \rangle$ is found to be $\zeta = 2.5$ MHz per photon. The mechanism is different from that of the asymmetric ATS observed in superconducting qubits with weak anharmonicity $[15,16,18]$. The extracted dip positions versus calibrated photon number $\langle n \rangle$ inside the resonator are presented in the inset. The solid lines show the calculated dip positions $\omega_n$ from the relations $\omega_n = \omega_{fg} \pm g_0\sqrt{\langle n \rangle} + 1 + \zeta \langle n \rangle$ including the correction of the ac Zeeman shift.

The simulated of the reflection coefficient is shown in Fig. 3(b) with a simplified model by numerically solving Eq. [2] in a steady state with the artificial atom parameters: $\omega_{eg} = 2\pi \times 3.379$ GHz, $\omega_{fg} = 2\pi \times 12.173$ GHz, $\gamma_{eg} = 2\pi \times 2$ MHz, $\gamma_{fg} = 2\pi \times 4$ MHz, $\gamma_{fe} = 2\pi \times 25$ MHz, $\gamma_{ff} = 2\pi \times 1$ MHz, and $\gamma_{ee} = 2\pi \times 20$ MHz. The dephasing effect comes from $|f\rangle \leftrightarrow |g\rangle$ fluctuations, which induced by the photon shot noise inside the cavity, is also added to $\gamma_{ff}$. The other second-order effects in the $|f\rangle \leftrightarrow |g\rangle$ interaction with the resonator are neglected. We simplify our model for the numerical calculations because of limited computer resources insufficient to operate with the full Hamiltonian in the dissipative regime with many photons because of the too large Hilbert space. Note that although there is a weak signature of zero-one state photon splitting, it cannot be resolved in the experiment because of the measurement noise. In our experiment, it is not yet reach the reliable resolution of the two states, which can be expressed as $\Delta\omega_{10} > 2\lambda_{fg}$ – the half-spacing between two dips is larger than $\lambda_{fg}$ which is the decay rate of the atomic off-diagonal terms $\rho_{fg}$, where $\Delta\omega_{10} = g_0(\sqrt{2} - 1) = 2\pi \times 27$ MHz and $\lambda_{fg} = 2\pi \times 16$ MHz (details are given in $[33]$).

Figure 3(c) shows the dependence of reflection power with $\langle n \rangle = 0$, $\langle n \rangle = 1$ and $\langle n \rangle = 5$. The bottom panel shows the difference between the traces with $\langle n \rangle = 0$ and $\langle n \rangle = 1$. It demonstrates the reflection of propagating microwaves in the transmission line could be controlled by microwave fields in the resonator at the single-photon level. It may find potential applications in single-photon switches, single-photon transistors, and so on. $[44,47]$

The intrinsic linewidth broadening with increasing driving power when $\langle n \rangle \gg 1$ is explained by the dephasing of $|f\rangle \leftrightarrow |g\rangle$ from the fluctuations of the photon inside the cavity $[19]$. This is contrast to the linewidth narrowing with increasing driving power in ATS driven with classical control fields $[17]$. The upper branch of the spectrum is broadened in the vicinity of 12.4 GHz, probably due to an unwanted low quality mode in the environment, which is difficult to remove at high (above 10 GHz) frequencies. Figure 3(d) shows the reflection power when $\langle n \rangle \approx 35$. The red line is fitted by a double-Gaussian curve of the form $A \exp(-2(\delta\omega/\Delta\omega)^2)$, where the fitting parameter $\Delta\omega/2\pi$ is found to be 121 MHz at 11.876 GHz and 186 MHz at 12.655 GHz. The linewidths are about twice the coupling strength $g_0$, which can be mainly determined by the Poisson photon distribution of coherent states formed inside the resonator $[33]$.

Finally, we change the coupling strength $g_0$ between the transition $|e\rangle \leftrightarrow |f\rangle$ and the resonator by tuning the magnetic field inside the device $\alpha$-loop $[34]$. We confirm that the splitting in the reflection spectrum of the coupled system around the transition frequency $\omega_{fe}$ is also twice the coupling strength $g_0$.

In conclusion, we engineer an interaction between a superconducting artificial atom and a resonator and observe the vacuum-induced (quantized) Autler-Townes
doublet by direct measurement of the reflection spectrum from the atom. Furthermore, we study the transition of the quantum-to-classical Autler-Townes doublet with changing the number of photons inside the resonator. Our work is an important step towards the control of propagating microwaves in a transmission line by a single photon confined in a different frequency resonator. With the further improvements, the effect may have potential applications in single-photon switches, single-photon transistors, and so on.

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