Single photons and unconventional photon blockade in quantum dot cavity-QED

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We observe the unconventional photon blockade effect in quantum dot cavity QED, which, in contrast to conventional photon blockade, operates in the weak coupling regime. A single quantum dot transition is simultaneously coupled to two orthogonally polarized optical cavity modes, and by careful tuning of the input and output state of polarization, the unconventional photon blockade effect is observed. We find a minimum second-order correlation \(g^{(2)}(0)\) \(\approx 0.37\) which corresponds to \(g^{(2)}(0) \approx 0.005\) when corrected for detector jitter, and observe the expected polarization dependency and photon bunching and anti-bunching very close-by in parameter space, which indicates the abrupt change from phase to amplitude squeezing.

A two-level system strongly coupled to a cavity results in a photon-number dependent energy. This dressing gives rise to the photon blockade effect [1, 2] resulting in photon-number dependent transmission and reflection, enabling the transformation of incident coherent light into specific photon number states such as single photons. Single-photon sources are a crucial ingredient for various photonic quantum technologies ranging from quantum key distribution to optical quantum computing. Such sources are characterized by a vanishing second-order auto-correlation \(g^{(2)}(0)\) \(\approx 0\) [3].

In the strong coupling regime, where the coupling between the two-level system and the cavity is larger than the cavity decay rate \((g > \kappa)\) [4], photon blockade has been demonstrated in atomic systems [5], quantum dots in photonic crystal cavities [6], and circuit QED [7, 8]. At the onset of the weak coupling regime \((g \approx \kappa)\), it has been shown that by detuning the dipole transition frequency with respect to the cavity resonance, photon blockade can still be observed [9]. However, moving further into the weak coupling regime \((g < \kappa)\) which is much easier to achieve, conventional photon blockade is no longer possible because the energy gap between the polariton states disappears. Nevertheless, also in the weak coupling regime, the two-level system enables photon number sensitivity, which has recently enabled high-quality single photon sources using polarization postselection [10–12] or optimized cavity in-coupling [13][14].

In 2010, Liew and Savona introduced the concept of the unconventional photon blockade effect [16] which operates even for weak non-linearities. It is a quantum interference effect between different excitation pathways which requires at least two or more degrees of freedom [17]. It was first investigated for Kerr non-linearities [16, 18], then for \(\chi^{(3)}\) non-linearities [19] and the Jaynes-Cummings [15] system which we focus on here. Both the conventional and unconventional photon blockade effect result in transmitted light with vanishing photon auto-correlation \(g^{(2)}(0) < 10^{-2}\) [18, 20], however, the underlying physical mechanisms are completely different, see Fig. 1. In the strong coupling regime, the nonlinearity of the dressed spectrum prevents reaching the two photon state for a particular laser frequency [red arrows in Fig. 1(a)], which is impossible in the weak coupling regime (green arrows). In unconventional photon blockade [Fig. 1(b)], reduction of the two-photon state is achieved by de-
In this paper, we show experimental evidence of unconventional photon blockade (UPB) in quantum dot cavity-QED. The sample consists of a layer of self-assembled InAs/GaAs quantum dots embedded in a micropillar cavity (maximum Purcell factor $F_p = 11.2$) grown by molecular beam epitaxy [22]. The quantum dot layer is embedded in a P–I–N junction, separated by a 27 nm thick tunnel barrier from the electron reservoir to enable tuning of the quantum dot resonance frequency by the quantum-confined Stark effect. Due to the quantum dot fine-structure structure splitting, we need to consider only one quantum dot transition, which interacts with both the H and V cavity modes.

We model our system using a Jaynes-Cummings Hamiltonian in the rotating wave approximation with $g \ll \kappa$. The Hamiltonian for two cavity modes and one quantum dot transition driven by a continuous wave laser is written as

$$H = (\omega_L - \omega^V_c) \hat{a}^\dagger_V \hat{a}^\dagger_V + (\omega_L - \omega^H_c) \hat{a}^\dagger_H \hat{a}^\dagger_H + (\omega_L - \omega_{QD}) \hat{\sigma}^\dagger \hat{\sigma} + g \left( \hat{a} \hat{b}^\dagger + \hat{a}^\dagger \hat{b} \right) + \eta H \left( \hat{a} H + \hat{a}^\dagger_H \right) + \eta V \left( \hat{a} V + \hat{a}^\dagger_V \right).$$

$\omega^H_c$ and $\omega^V_c$ are the resonance frequencies of the linearly polarized cavity modes, $\hat{a}^\dagger_H$ and $\hat{a}^\dagger_V$, the photon creation operators, $\omega_{QD}$ is the quantum dot resonance frequency, and $\hat{\sigma}^\dagger$ the exciton creation operator. $\hat{b} = \hat{a}_V \cos \phi + \hat{a}_H \sin \phi$ is the cavity photon annihilation operator along the quantum dot dipole orientation, and $\phi$ is the relative angle. In our case the angle is $\phi = 94^\circ$, which means that the H-cavity mode couples better to the exciton transition. $\eta H$ and $\eta V$ are the amplitudes of the incident coherent light coupling to the H and V cavity modes. For numerical simulations, we add relaxation of the cavity modes and dephasing of the quantum dot transition and solve the corresponding quantum master equation [12, 23–25], add the output polarizer and calculate the mean photon number and second order correlation function. All theoretically obtained $g^{(2)}(\tau)$ data is convoluted with the detector response (530 ps) to match the experimental setup.

Fig. 2 shows how the second order correlation $g^{(2)}(\tau = 0)$ of the transmitted photons depends on the linear input and linear output polarization angle. In all current quantum dot based single photon sources [10–12], only one cavity mode is excited with the laser, and by using a crossed polarizer, single photons are obtained. This condition is indicated with arrow A in Fig. 2. We see that, by using an input polarization where both polarization modes are excited (indicated by arrow B), a much larger range is found where single photons can be produced. This is where the unconventional photon blockade can be observed.

Now, we investigate more closely region B of Fig. 2, where both cavity modes are excited ($\theta_{in} = 45^\circ$). Furthermore, we add the experimentally unavoidable polarization splitting of the H and V cavity modes which is 10 GHz for the device under investigation. Furthermore, we vary the detected output polarization in the most general way, by introducing $\lambda/2$ and $\lambda/4$ wave plates before the final polarizer in the transmission path. The experimental setup is sketched in the inset of Fig. 3(b). Fig. 3(b) shows how this polarization projection affects the mean photon number $\langle n_{out} \rangle$, for $\langle n_{in} \rangle = (n_h + n_v)^2 = 0.06$ in the simulation and in the experiment [Fig. 3(a)]. This
region is highly dependent on the cavity splitting and the quantum dot dipole angle, careful determination of the parameters allows us to obtain good agreement to experimental data [Fig. 3(a)]. In this low mean photon number region, the second-order correlation \( g^{(2)}(0) \) shows a non-trivial behavior as a function of the output polarization state, shown in Fig. 3(c, experiment) and (d, theory): First, we observe the expected unconventional photon blockade anti-bunching (blue region). The experimentally measured minimum \( g^{(2)}(0) \) is 0.37, which is limited by the detector response function. The theoretical data which takes the detector response into account agrees very well to the experimental data and predicts a bare \( g^{(2)}(0) \approx 0.005 \). Second, we find that, close-by in parameter space, there is a region where bunched photons are produced. This enhancement of the two-photon probability happens via constructive interference leading to phase squeezing. Theoretical and experimental data show good agreement, we attribute the somewhat more extended antibunching region to long-time drifts during the course of the experiment (10 hours).

In Fig. 3(e) and 3(f) we show the two-time correlation function \( g^{(2)}(\tau) \) for the two cases indicated by the arrows. The observed width and height of the anti-bunching and bunching peak predicted by the theory is in excellent agreement with the observed experimental data. For two coupled Kerr resonators in the UPB regime, one observes oscillations in \( g^{(2)}(\tau) \) when collecting the output of only one of the cavities [16]. During finalizing this paper, a manuscript describing a first observation of this effect has appeared [26]. In our case, these oscillations are absent because the system works mostly as a unidirectional dissipative coupler [27], and the photon field behind the output polarizer contains contributions from both cavities modes, which suppresses the oscillations in \( g^{(2)}(\tau) \).

An alternative way to understand the unconventional photon blockade is in terms of Gaussian squeezed states [28]: For any coherent state \( |\alpha\rangle \), there exists an optimal squeeze parameter \( \xi \) that minimizes the two-photon correlation \( g^{(2)}(0) \), which can be made vanishing for a weak driving fields. We find that, even with a small amount of squeezing, it is possible to significantly reduce the two-photon distribution and minimize \( g^{(2)}(0) \) for low mean photon numbers. A Gaussian squeezed state is produced from vacuum like \( D(\alpha)S(\xi)|0\rangle = |\alpha, \xi\rangle \). Here \( S \) is the squeezing operator with \( \xi = r \exp^{i\theta} (0 \leq r < \infty, 0 \leq \theta \leq 2\pi) \). \( D \) is the displacement operator, and the complex displacement amplitude \( \alpha = \bar{\alpha} \exp^{i\vartheta} (0 \leq \bar{\alpha} < \infty, 0 \leq \vartheta \leq 2\pi) \). For \( \theta = \vartheta = 0 \), we can calculate the two photon probability in the small-\( \alpha \) (low mean photon number) limit as

\[
|2D(\alpha)S(\xi)|0\rangle|^2 \approx (\bar{\alpha}^2 - r^2)/2, \tag{1}
\]

using a Taylor expansion. We see that, in order to obtain a vanishing two-photon probability, the squeeze parameter \( r \) needs to be equal to \( \bar{\alpha}^2 \) which is the mean photon number. By defining the amount of quadrature squeezing as \( \langle (\Delta X)^2 \rangle = \frac{1}{2} e^{-2r} \) and considering a \( \langle n_{out} \rangle \approx 0.004 \) (Fig. 3(a)), this condition leads to \( 10 \log_{10}(e^{-0.008}) = -3 \times 10^{-2} \) dB squeezing. Interestingly, this result means that, for a weak coherent state, only a very small amount of squeezing is needed to make \( g^{(2)}(0) \) drop to zero.

In Fig. 4 we show further analysis of the theoretical calculations for the experimental state produced by the unconventional photon blockade as indicated by arrow D in Fig. 2(c) and (d). In agreement with equation (1) we observe that the 2-photon state in the photon number distribution shown in Fig. 4(a) is suppressed. Consequently, from the photon number variance given in Fig. 4(b), we observe that the state is amplitude squeezed. Further, by moving from the region of arrow C to D in Fig. 3(c), the observed state switches from a phase squeezed to an amplitude squeezed state, which is a clear signature of the unconventional photon blockade effect [17].

Finally, we discuss whether the UPB effect can be used to enhance the performance of single photon sources, and...
Figure 4. (a) Calculated photon number distribution of a coherent state and for the condition indicated by the arrow C and D in Fig. 2(c). (b) The calculated photon number variance for the states presented in (a) showing amplitude squeezing in the region where we observe the unconventional photon blockade. (c) Mean photon number \( \langle n_{\text{out}} \rangle \) as a function of input polarization. We see that a large improvement of the single photon brightness can be obtained by exploiting the UPB effect. The simulation is performed for three cavity splittings \( \Delta f_{\text{cav}} \) showing that the enhancement is largest in a polarization degenerate cavity.

in particular their brightness. Traditionally, the quantum dot is excited by one linearly polarized cavity mode and photons are collected via the orthogonal mode. In our experiment, the quantum dot excitation probability is 1 − \( \cos(45^\circ) \approx 0.0024 \), and, once excited, it has 1 − 0.0024 chance to emit into the collection cavity mode, which leads to a low total efficiency. In the unconventional photon blockade regime, arrow B in Fig 2, this efficiency is higher. To further explore this, we show in Fig. 4(c) the mean photon number \( \langle n_{\text{out}} \rangle \) as a function of the input polarization with constant input laser power (the polarization output state is chosen such that \( g^{(2)}(0) \approx 0 \)). We see that, by rotating the input polarization from 0° to 45°, the output mean photon number can be increased by approximately a factor 10. The simulation is done for various cavity splittings \( \Delta f_{\text{cav}} \) which shows that increasing the cavity splitting reduces this enhancement.

In conclusion, we have experimentally observed the unconventional photon blockade effect using a single quantum dot resonance coupled to two orthogonally polarized cavity modes. We find the expected drop in \( g^{(2)}(0) \), but additionally and very close in parameter space, we also find that the transmitted light statistics can be tuned from anti-bunched to bunched, all in good agreement to theoretical models and simulations. In contrast to conventional photon blockade, no energy splitting of the polariton resonances is required, allowing to obtain \( g^{(2)}(0) \approx 0 \) even with weak non-linearities. Finally, under certain conditions, we find that the unconventional photon blockade effect can increase the brightness of the single photon sources.

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