Phonon Mediated Off-Resonant Quantum Dot-Cavity Coupling

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A theoretical model for the phonon-mediated off-resonant coupling between a quantum dot and a cavity, under resonant excitation of the quantum dot, is presented. We show that the coupling is caused by electron-phonon interaction in the quantum dot and is enhanced by the cavity. We analyze recently observed resonant quantum dot spectroscopic data by our theoretical model.

One of the interesting recent developments in cavity quantum electrodynamics (CQED) experiments with quantum dots (QDs) coupled to semiconductor microcavities is the observation of off-resonant dot-cavity coupling. This unusual phenomenon is observed both in photoluminescence studies, where an above-band laser generates carriers that incoherently relax to recombine at the QD frequency [1–3] and also under resonant excitation of the QD or the cavity [4, 5]. The coupling observed via photoluminescence is attributed to several phenomena including pure QD dephasing [6], the electron-phonon interaction [7, 8], multi-exciton complexes [9] and charges in the vicinity of the QD [10]. To isolate the role of phonons in off-resonant QD-cavity coupling, studies employing resonant excitation of the QD are preferable as they avoid possible complications arising from multi-excitonic complexes and nearby charges generated in above band pumping. Resonant excitation of a QD coupled to an off-resonant cavity has been recently used to perform QD spectroscopy, enabling the observation of power broadening of the QD line-width and saturation of the cavity emission [11–13]. However, there is presently no theoretical model accounting for off-resonant dot-cavity coupling for the case of resonant QD excitation. Although off-resonant coupling can be modeled by introducing a phenomenological incoherent cavity pumping mechanism [14], such a treatment masks the actual physical phenomenon responsible for the coupling. Without an explicit coherent driving term in the system Hamiltonian, the resonant QD spectroscopy results cannot be explained.

In this Letter, we theoretically model the phonon-mediated interaction between a cavity mode and an off-resonant QD under resonant excitation of the QD. We first model the coupling via pure QD dephasing. Then, we propose a new model where the phonon-mediated coupling is enhanced by the presence of the cavity. We compare these two models and find that they provide qualitatively similar signatures in terms of experimental resonant QD spectroscopic studies, such as power broadening and saturation of a resonantly driven QD [11, 12]. However, in the newly proposed model, the coupling is maintained at very large QD-cavity detunings (∼ 3 meV), as observed in the recent experiments [11]. We also observe the signature of the inherent asymmetry between phonon emission and absorption rate in the simulated QD spectroscopy results, depending on whether the QD is red or blue detuned from the cavity.

The dynamics of a coherently driven QD (with a ground state |g⟩ and an excited state |e⟩) coupled to an off-resonant cavity mode is governed by the Hamiltonian

\[ H = \Delta \omega_c a^\dagger a + \Delta \omega_a \sigma^\dagger \sigma + i g (a^\dagger \sigma - a \sigma^\dagger) + \Omega (\sigma + \sigma^\dagger) \] (1)

where, \( a \) and \( \sigma \) are the annihilation and lowering operators for the cavity mode and the QD, respectively; \( \Delta \omega_c = \omega_c - \omega_l \) and \( \Delta \omega_a = \omega_a - \omega_l \) are the cavity and dot detunings from the driving laser, respectively; \( \Delta = \omega_a - \omega_c \) is the QD-cavity detuning; \( \Omega \) is the Rabi frequency of the driving laser and is proportional to the laser field amplitude and \( g \) is the coherent interaction strength between the QD and the cavity.

In this coupled system, there are two independent mechanisms for energy dissipation: cavity decay and QD dipole decay. The system losses can be modeled by the Liouvillian, and the Master equation describing the dynamics of the lossy system is given by

\[ \frac{d \rho}{dt} = -i [H, \rho] + 2 \kappa \mathcal{L}[a] + 2 \gamma \mathcal{L}[\sigma] \] (2)

where, \( \rho \) is the density matrix of the coupled QD-cavity system, \( 2 \gamma \) and \( 2 \kappa \) are the QD spontaneous emission rate and the cavity population decay rate, respectively. We neglect any non-radiative decay of the QD exciton. \( \mathcal{L}[D] \) is the Lindblad operator corresponding to a collapse operator \( D \) and is given by:

\[ \mathcal{L}[D] = D \rho D^\dagger - \frac{1}{2} D^\dagger D \rho - \frac{1}{2} \rho D^\dagger D \] (3)

In addition, phonons in the solid state system destroy the coherence of the exciton. This is generally modeled by adding an additional incoherent decay term \( 2 \gamma_d \mathcal{L}[\sigma^\dagger \sigma] \) to the Master equation, where \( 2 \gamma_d \) is the pure dephasing rate of the QD. This term destroys the polarization of the QD without affecting the population of the QD. The dissipation of the QD polarization and population \( \sigma_z = \)
\[\langle \sigma^1, \sigma \rangle \] is given by the following mean-field equations:

\[
\frac{d\langle \sigma \rangle}{dt} = -(\gamma + \gamma_d)\langle \sigma \rangle
\]  
(4)

\[
\frac{d\langle \sigma_z \rangle}{dt} = -2\gamma(1 + \langle \sigma_z \rangle)
\]  
(5)

Hence, the linewidth of the QD, at the zero excitation power limit, is given by \(2(\gamma + \gamma_d)\). However, in this model, the effect of phonons is embedded in the phenomenological pure dephasing rate \(\gamma_d\), which affects only the QD, and does not include any cavity effects.

We now propose a different model for off-resonant dot-cavity coupling, where the phonon-mediated coupling strength is affected by both the cavity and the QD. The effect of phonons can be modeled by replacing the pure dephasing term with two additional incoherent decay terms in the Master equation. For blue-detuned QD (Fig. 1 a) we have:

\[
\frac{d\rho}{dt} = -i[H, \rho] + 2\kappa L[a] + 2\gamma L[\sigma] \\
+ 2\gamma_r n L(\sigma^1 a) + 2\gamma_r(n + 1) L(\sigma a^1)
\]  
(6)

where, \(2\gamma_r\) is an effective decay rate of the QD exciton states via the emission of a phonon and a photon at the off-resonant cavity frequency. \(n\) is the average number of phonons at the dot-cavity detuning frequency \((\Delta)\) present in the system at thermal equilibrium with the reservoir at a temperature \(T\), and is given by

\[
n(\Delta, T) = \frac{1}{e^{h\Delta/k_B T} - 1}
\]  
(7)

The analysis for a QD red detuned from the cavity (Fig. 1 b) can be carried out in a similar manner by replacing the final two terms of Eqn.6 with \(2\gamma_r n L(\sigma a^1)\) and \(2\gamma_r(n + 1) L(\sigma^1 a)\).

The decay term \(\sigma^1 a\) denotes the annihilation of a cavity photon and excitation of the QD, while the term \(\sigma a^1\) denotes the creation of a cavity photon and collapse of the QD to its ground state, accompanied by the creation (or annihilation) of phonons to compensate for the QD-cavity frequency difference. Only the second process is important for observing cavity emission under resonant excitation of the dot. We also note that the observation of QD emission under resonant excitation of the cavity [4] can be modeled in the same way by changing the coherent driving term from \(\Omega(\sigma + \sigma^1)\) to \(\Omega(a + a^1)\). In this situation, the collapse operator \(\sigma^1 a\) is important. Similar decay channels have been proposed to model cavity assisted atomic decay [15]. A detailed derivation of these incoherent terms can be found in the Supplementary Materials. We call this process a cavity-enhanced phonon process.

We first consider the case where the QD is blue detuned from the cavity (Fig. 1 a). The qualitative nature of the dissipation of the QD polarization and population can be determined by the mean-field equations (assuming \(n = 0\), i.e., at the zero temperature limit):

\[
\frac{d\langle \sigma \rangle}{dt} = -\gamma \langle \sigma \rangle - \gamma_r (1 + \langle a^1 a \rangle) \langle \sigma \rangle
\]  
(8)

\[
\frac{d\langle \sigma_z \rangle}{dt} = -2\gamma(1 + \langle \sigma_z \rangle) - 2\gamma_r(1 + \langle a^1 a \rangle)(1 + \langle \sigma_z \rangle)
\]  
(9)
We notice that, unlike the pure dephasing case, both the QD population and polarization are affected by the cavity enhanced phonon process. The linewidth of the QD at the zero excitation power limit is given by \( \Delta \omega_0 = \frac{1}{\gamma_0} \). We observe that at large detuning, \( \Delta \omega \) shows the cavity fluorescence as a function of the laser Rabi frequency \( \Omega \). For both models, the cavity fluorescence \( I \) follows a saturation curve

\[
I = I_{\text{sat}} \frac{\hat{P}}{1 + \hat{P}}
\]

where, \( \hat{P} \propto \Omega^2 \) and \( I_{\text{sat}} \) is the saturated cavity emission intensity. As noted previously in the article, the cavity emission is higher when the process is enhanced by the presence of the cavity. Furthermore, we investigate the dependence of the saturation emission intensity \( I_{\text{sat}} \) as a function of the QD-cavity detuning \( \Delta \) (Fig. 2 c). \( I_{\text{sat}} \) falls of as \( 1/\Delta^2 \) with the detuning \( \Delta \) when the dot-cavity coupling is modeled as a pure dephasing process. However, when the coupling is modeled as a cavity enhanced phonon process, the saturation intensity exhibits a diminished dependence on detuning \( \Delta \) (estimated to be \( \sim 1/\Delta^{0.25} \)). Hence, one may observe off-resonant coupling for larger detunings when the phonon process is enhanced by the cavity. Fig. 2 d shows \( \log(I) \) as a function of the cavity decay rate \( \kappa \), for a fixed detuning of \( \Delta/2\pi = 200 \text{ GHz} \). For both models, the emission falls off as \( 1/\kappa^2 \), signifying that the off-resonant coupling does not depend on the overlap between the QD and the cavity spectra.

We now measure the QD line-width \( \Delta \omega \) monitoring the cavity emission, while scanning the laser wavelength across the QD resonance, similar to the experiments in Ref. [11]. We observe that at very low excitation power \( \Omega/2\pi = 1 \text{ GHz} \) and large QD-cavity detuning \( \Delta/2\pi = 12\kappa \), the linewidths of the QD are very close to the theoretical linewidth in the absence of the cavity (shown by the solid black line in Fig. 3 a). At a constant QD-cavity detuning and laser excitation power, the QD line-width increases with increasing \( \gamma_0 \) and \( \gamma_r \) (Fig. 3 a). We observe broadening of the line-width with increasing laser power (Figs. 3 b). The power broadened QD linewidth \( \Delta \omega \) is fit with the model \( \Delta \omega = \Delta \omega_0 \sqrt{1 + \hat{P}} \), where \( \Delta \omega_0 \) is the intrinsic line-width of the QD and \( \hat{P} \) is obtained from the fit to the saturation of the cavity emission [11]. The theoretical model does not reproduce the additional power-independent broadening of the QD [11] and we believe that this extra broadening may result from QD spectral diffusion [19]. We analyze the intrinsic QD linewidth \( \Delta \omega_0 \) (without power broadening, i.e., obtained from plots in Figs. 3 b at \( \Omega = 0 \) limit) as a function of the dot-cavity detuning \( \Delta \) for two different models (Figure 3 c). We observe that at large detuning, \( \Delta \omega_0 \) approaches the unperturbed QD linewidth \( 2(\gamma + \gamma_d) \) and \( 2(\gamma + \gamma_c) \), respectively. We fit empirical models of \( \Delta^{-\alpha} \) to the intrinsic linewidths for the two models and find that with pure QD dephasing, \( \Delta \omega_0 \) falls off more slowly (\( \alpha \approx 0.4 \)) compared to the cavity enhanced coupling (\( \alpha \approx 0.7 \)). The weak dependence of the intrinsic QD linewidths on the dot-cavity detuning shows that the off-resonant cavity does not perturb the QD significantly.
However, the results are dramatically different for a red-detuned QD. This is a result of the fact that the rate of the incoherent process involving the operator $\sigma a \dagger$ is different depending on whether the QD is red or blue detuned from the cavity (because at any temperature, the rates of absorption and emission of phonons are different). This asymmetry is manifested in both the emission collected at the cavity resonance and in the QD spectroscopy result. Fig. 4 a shows the difference in the QD linewidths as a function of the bath temperature $T$ for different driving laser Rabi frequencies $\Omega$, when the QD is red or blue detuned from the cavity by same absolute value. We observe that the difference in linewidth is higher when the QD is weakly driven and hence is not power broadened. Full quantum optical simulations show that the linewidth depends on the bath temperature, but reaches the maximum value of $2\gamma_\rho$ at a higher temperature. Fig. 4 b shows the ratio of the cavity emission as a function of the bath temperature for different driving laser Rabi frequencies $\Omega$. The difference in cavity intensity is maximum at lower bath temperature and is almost zero at higher temperature.

In conclusion, we have proposed a theoretical model for the off-resonant dot-cavity coupling under resonant excitation of the QD. Introduction of an incoherent decay channel (referred to as the cavity enhanced phonon process, which is different from pure dephasing) shows that the phonon-mediated dot-cavity coupling is enhanced by the presence of the cavity. By comparing the power-broadening and saturation of the QD between the cavity enhanced phonon process and pure QD dephasing, we found that the dot-cavity coupling is enhanced in the former case and is observed for a larger QD-cavity detuning. Our model can also be used to explain asymmetry in the spectroscopy results for a QD blue and red detuned from the cavity. We believe that such an off-resonant dot-cavity coupling can be used as an efficient read-out channel for resonant QD spectroscopy and for QD-spin manipulation.

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I. SUPPLEMENTARY MATERIALS

A. Derivation of Decay Terms

We use the level diagram as shown in Fig. 1 a to model the effect of phonons explicitly. The Hamiltonian of the system is given by

\[ H = H_0 + H_I \]  

(12)

where,

\[ H_0 = \omega_1 \left| 1 \right> \left< 1 \right| + \omega_2 \left| 2 \right> \left< 2 \right| + \omega_3 \left| 3 \right> \left< 3 \right| + \omega a^\dagger a + \sum_j \nu_j b_j^\dagger b_j \]  

(13)

and

\[ H_I = g_\nu (a \left| 3 \right> \left< 1 \right| + a^\dagger \left| 1 \right> \left< 3 \right|) + \sum_j g_{23}^j (b_j \left| 3 \right> \left< 2 \right| + b_j^\dagger \left| 2 \right> \left< 3 \right|) \]  

(14)

where, \( \left| i \right> \left< i \right| \) is the population operator for \( i^{th} \) level; \( a \) is the annihilation operator for the cavity mode; \( b_j \) is the annihilation operator for a phonon in the \( j^{th} \) mode. \( \omega_i \), \( \omega \) and \( \nu_j \) are the frequencies of the \( i^{th} \) energy level, cavity resonance and a phonon in the \( j^{th} \) mode. \( g_\nu \) signifies the interaction strength between the cavity and the virtual transition and \( g_{23}^j \) is the interaction strength between the QD exciton and an \( j^{th} \) mode phonon. We note that this interaction Hamiltonian is valid only for the level structure as in Fig. 1 a, where the cavity is at lower energy than the QD. For the situation in Fig. 1 b (where the cavity is of higher energy compared to the QD), the interaction Hamiltonian \( H_I \) is given by

\[ H_I = g_\nu (a \left| 3 \right> \left< 1 \right| + a^\dagger \left| 1 \right> \left< 3 \right|) + \sum_j g_{23}^j (b_j \left| 3 \right> \left< 2 \right| + b_j^\dagger \left| 2 \right> \left< 3 \right|) \]  

(15)

In the following derivation, we will use the situation shown in Fig. 1 a.

If we define the QD resonance frequency as \( \omega_a \), then \( \omega_a = \omega_2 - \omega_1 \); and the cavity frequency is given by \( \omega_c = \omega_3 - \omega_1 \). Then the QD-cavity detuning is given by \( \Delta = \omega_2 - \omega_3 \). Defining \( \sigma_{ij} = |i\rangle \langle j| \), we can write

\[ \sigma_{13} = -i[\sigma_{13}, H_0 + H_I] \]

\[ = -i\omega_c \sigma_{13} - ig_\nu a (\sigma_{11} - \sigma_{33}) - i \sum_j g_{23}^j b_j^\dagger \sigma_{12} \]

Similarly,

\[ \sigma_{23} = -i[\sigma_{23}, H_0 + H_I] \]

\[ = i \Delta \sigma_{23} - ig_\nu a \sigma_{21} - i \sum_j g_{23}^j b_j^\dagger (\sigma_{22} - \sigma_{33}) \]

Separating the slow and the fast components of the operators, we can write

\[ \dot{\sigma}_{13} = -i g_\nu a (\sigma_{11} - \sigma_{33}) - i \sum_j g_{23}^j b_j^\dagger \sigma_{12} \]

(21)

and

\[ \dot{\sigma}_{23} = i (\Delta - \omega_j) \sigma_{23} - ig_\nu a \sigma_{21} - i \sum_j g_{23}^j b_j^\dagger (\sigma_{22} - \sigma_{33}) \]

(22)

As level 3 is a virtual level, it is never populated. Hence by adiabatic elimination, using \( \dot{\sigma}_{13} = \dot{\sigma}_{23} = 0 \), we obtain

\[ \sigma_{23} = \frac{g_\nu a \sigma_{21} + \sum_j g_{23}^j b_j^\dagger (\sigma_{22} - \sigma_{33})}{\Delta - \omega_j} \]

(23)

and

\[ \dot{\sigma}_{12} = -g_\nu a (\sigma_{11} - \sigma_{33}) \]

(24)
Using these values, we can find the interaction Hamiltonian. The first term $g_{j}^{d}(\sigma_{31} + a^{\dagger}\sigma_{13})$ denotes the coherent dynamics. The second term, which signifies the effect of phonons, can be written as (using the Eqns. 23 and 24)

$$H_{ph} = \sum_{j} g_{j}^{23}(b_{j}^{\dagger}\sigma_{32} + b_{j}\sigma_{23})$$

$$= \sum_{j} g_{j}^{23}(b_{j}^{\dagger}\bar{\sigma}_{32} + \bar{b}_{j}\bar{\sigma}_{23})$$

$$= \sum_{j} \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}}(b_{j}^{\dagger}\bar{\sigma}_{12} + \bar{b}_{j}\bar{\sigma}_{21})$$

$$+ \sum_{j} \frac{(g_{j}^{23})^{2}}{\Delta - \omega_{j}}(\bar{\sigma}_{22} - \bar{\sigma}_{33})(b_{j}^{\dagger}\bar{b}_{j} + \bar{b}_{j}b_{j})$$

The second term involves two-phonon processes which are less likely. If we neglect them, we can model the effect of phonons as follows:

$$H_{ph} = \sum_{j} \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}}(b_{j}^{\dagger}\bar{\sigma}_{12} + \bar{b}_{j}\bar{\sigma}_{21})$$

We can write this Hamiltonian as

$$H_{ph} = (\bar{\sigma}_{12}\bar{\Gamma}^{\dagger} + \bar{a}\bar{\sigma}_{21}\bar{\Gamma})$$

where, the operator $\bar{\Gamma}$ can be written as

$$\bar{\Gamma} = \sum_{j} \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}}\bar{b}_{j}$$

and

$$\Gamma = \sum_{j} \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}}b_{j}e^{-i\omega_{j}t}$$

To obtain the familiar Lindblad term, we take the partial trace of the correlation between the reservoir operators over the reservoir variables. The correlation is given by

$$\langle \Gamma^{\dagger}(t')\Gamma(t) \rangle = \sum_{j} \left| \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}} \right|^{2} e^{-i\omega_{j}(t-t')}\langle b_{j}^{\dagger}b_{j} \rangle_{R}$$

As the operators $b_{j}$ are bosonic and the system is in thermal equilibrium with a bath at temperature $T$, using the relation

$$\langle b_{j}^{\dagger}b_{j} \rangle_{R} = \bar{n}(\omega_{j}, T) = \frac{1}{e^{\hbar\omega_{j}/k_{B}T} - 1}$$

we find

$$\langle \Gamma^{\dagger}(t')\Gamma(t) \rangle = \sum_{j} \left| \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}} \right|^{2} e^{-i\omega_{j}(t-t')}\bar{n}(\omega_{j}, T)$$

and

$$\langle \Gamma(t')\Gamma^{\dagger}(t) \rangle_{R} = \sum_{j} \left| \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}} \right|^{2} e^{-i\omega_{j}(t-t')}\bar{n}(\omega_{j}, T) + 1$$

From the correlation, we find that the phonons with frequency $\Delta$ (corresponding to the difference between levels $2$ and $3$), i.e., QD-cavity mode detuning, have the maximum contribution in the interaction Hamiltonian. In the Born-Markov approximation, we can model the electron-phonon interaction (for Fig. 1 a) as an incoherent decay process by adding two extra terms to the Master equation: $2\gamma_{r}\bar{n}\mathcal{L}(\sigma^{\dagger}a)$ and $2\gamma_{r}(\bar{n} + 1)\mathcal{L}(\sigma^{\dagger}a)$, $\gamma_{r}$ being the effective decay rate of the excited QD state and is given by

$$\gamma_{r} = \frac{1}{2} \sum_{j} \left| \frac{g_{j}^{23}g_{v}}{\Delta - \omega_{j}} \right|^{2}$$

For the situation shown in Fig. 1 b, the decay terms are given by $2\gamma_{r}\bar{n}\mathcal{L}(\sigma^{\dagger}a)$ and $2\gamma_{r}(\bar{n} + 1)\mathcal{L}(\sigma^{\dagger}a)$. We note that the different rates in both cases are due to an inherent asymmetry between the absorption and emission rates of the phonons.

**B. Derivation of the Mean Field Equations**

To find the mean field equations for an operator $A$ from the Master equation, we used the following relation:

$$\frac{d \langle A \rangle}{dt} = \frac{d}{dt} Tr[A \rho] = Tr \left[ A \frac{d \rho}{dt} \right]$$

For the cavity enhanced phonon process, the mean field equations for a non-zero ($\bar{n}$) is given by (when the QD is blue detuned from the cavity)

$$\frac{d \langle \sigma \rangle}{dt} = -\gamma_{r}\langle \sigma \rangle - \gamma_{r}(1 + \langle a^{\dagger}a \rangle)\langle \sigma \rangle$$

$$-2\gamma_{r}(1 + \langle \sigma_{z} \rangle)$$

When the QD is red detuned from the cavity, the mean field equations are:

$$\frac{d \langle \sigma \rangle}{dt} = -\gamma_{r}\langle \sigma \rangle - \gamma_{r}(1 + \langle a^{\dagger}a \rangle)\langle \sigma \rangle$$

$$-2\gamma_{r}(1 + \langle \sigma_{z} \rangle)$$

We note that while deriving these mean-field equations, we assume that the cavity and QD operators are uncorrelated and write

$$\langle a^{\dagger}a\sigma \rangle = \langle a^{\dagger} \rangle \langle \sigma \rangle$$
Under weak driving, the approximation holds very well. However, in full quantum optical simulations, we do not make any assumptions.

[1] K. Hennessy, A. Badolato, M. Winger, D. Gerace, M. Atature, S. Gulde, S. Falt, E. L. Hu, and A. Imamoglu, Nature 445, 896 (2007).

[2] D. Press, S. Götzinger, S. Reitzenstein, C. Hofmann, A. Löffler, M. Kamp, A. Forchel, and Y. Yamamoto, Phys. Rev. Lett. 98, 117402 (2007), URL http://link.aps.org/abstract/PRL/v98/e117402.

[3] M. Kaniber, A. Laucht, A. Neumann, J. M. Villas-Boas, M. Bichler, M.-C. Amann, and J. J. Finley, Physical Review B (Condensed Matter and Materials Physics) 77, 161303 (2008), URL http://link.aps.org/abstract/PRB/v77/e161303.

[4] D. Englund, A. Majumdar, A. Faraon, M. Toishi, N. Stoltz, P. Petroff, and J. Vučković, Phys. Rev. Lett. 104, 073904 (2010).

[5] S. Ates, S. M. Ulrich, A. Ulhaq, S. Reitzenstein, A. Löffler, S. Höfling, A. Forchel, and P. Michler, Nature Photonics 3, 724 (2009).

[6] A. Auffeves, J.-M. Gerard, and J.-P. Poizat, Phys. Rev. A 79, 053838 (2009), URL http://link.aps.org/abstract/PRA/v79/e053838.

[7] Y. Ota, S. Iwamoto, N. Kumagai, and Y. Arakawa, arXiv:0908.0788v1 [cond-mat.mes-hall] (2009).

[8] U. Hohenester, Phys. Rev. B 81, 155303 (2010).

[9] M. Winger, T. Volz, G. Tarel, S. Portolan, A. Badolato, K. J. Hennessy, E. L. Hu, A. Beveratos, J. Finley, V. Savona, et al., Phys. Rev. Lett. 103, 207403 (2009), URL http://link.aps.org/abstract/PRL/v103/e207403.

[10] N. Chauvin, C. Zinoni, M. Francardi, A. Gerardino, L. Balet, B. Alloing, L. H. Li, and A. Fiore, Physical Review B (Condensed Matter and Materials Physics) 80, 241306 (2009), URL http://link.aps.org/abstract/PRB/v80/e241306.

[11] A. Majumdar, A. Faraon, E. D. Kim, D. Englund, H. Kim, P. Petroff, and J. Vučković, Phys. Rev. B 82, 045306 (2010).

[12] A. Ulhaq, S. Ates, S. Weiler, S. M. Ulrich, S. Reitzenstein, A. Löffler, S. Höfling, L. Worschech, A. Forchel, and P. Michler, Phys. Rev. B 82, 045307 (2010).

[13] E. D. Kim, A. Majumdar, H. Kim, P. Petroff, and J. Vučković, Applied Physics Letters 97, 053111 (2010), URL http://link.aip.org/link/?APL/97/053111/1.

[14] F. P. Laussy, E. del Valle, and C. Tejedor, Phys. Rev. B 79, 235325 (2009).

[15] G. Gangopadhyay, S. Basu, and D. S. Ray, Phys. Rev. A 47, 1314 (1993).

[16] C. W. Gardiner and P. Zoller, Quantum Noise (Springer-Verlag, 2005).

[17] S. M. Tan, Journal of Optics B: Quantum and Semiclassical Optics 1, 424 (1999), URL http://stacks.iop.org/1464-4266/1/i=4/a=312.

[18] M. O. Scilly and M. S. Zubairy, Quantum Optics (Cambridge University Press, 2005).

[19] H. Kamada and T. Kutsuwa, Phys. Rev. B 78, 155324 (2008).