Feshbach blockade: Single-photon nonlinear optics using resonantly enhanced cavity polariton scattering from biexciton states

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Abstract – We theoretically demonstrate how the resonant coupling between a pair of cavity polaritons and a biexciton state can lead to a large single-photon Kerr nonlinearity in a semiconductor solid-state system. A fully analytical model of the scattering process between a pair of cavity polaritons is developed, which explicitly includes the biexcitonic intermediate state. A dramatic enhancement of the polariton-polariton interactions is predicted in the vicinity of the biexciton Feshbach resonance. Application to the generation of non-classical light from polariton dots is discussed.

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A number of quantum optical applications crucially rely on having a strong value of the effective photon-photon interaction mediated by the optical nonlinearity of the atomic or solid-state medium. The most interesting quantum physics appears in fact when the presence of a single photon is sufficient to significantly modulate the response of a device. As a simplest example, a stream of strongly antibunched photons is emitted by a cavity as soon as the nonlinear shift of the resonance by a single photon is larger than the cavity linewidth [1]. This physics is even more intriguing when more complex devices are considered: photons have been predicted to fermionize into Tonks-Girardeau gases as soon as the impenetrability condition is satisfied in a one-dimensional geometry. To this end, both strongly nonlinear optical fibers [2] and arrays of many coupled nonlinear cavities [3] have been considered. Quantum phase transitions between a coherent phase of “superfluid” photons to a Mott insulator one have also attracted a significant interest [4]. So far, the main obstacle against an experimental realization of all these phenomena has been the lack of scalable optical devices with sufficiently large nonlinearities and weak enough losses.

A great deal of the recent advances in the field of strongly correlated atomic gases have been made possible by the discovery of the so-called Feshbach resonance effect in atom-atom collisions [5]: the scattering cross-section is dramatically enhanced when the energy of the pair of colliding atoms is resonant with a long-lived quasi-bound molecular state. In typical experiments, an external magnetic field is used to tune the energy of the quasi-bound molecular state close to the energy zero of scattering states. At this point, the low-energy scattering amplitude diverges and a dilute ultra-cold atomic gas has the strongest possible interaction. Remarkable recent experiments with atomic Fermi gases have exploited such Feshbach resonances to demonstrate superfluid behavior at temperatures comparable to the Fermi temperature [6].

In this letter, we show how a similar Feshbach resonance mechanism can be used in a solid-state context to dramatically enhance the strength of the interactions between cavity polaritons in a planar microcavity in the strong-coupling regime [7]. During the last two decades, a significant literature has appeared on the physics of biexcitons in semiconductor materials, i.e. two electron-two hole bound complexes analogous to hydrogen molecules [8]. A recent work by Wouters [9] has pointed out the possibility of exploiting biexciton Feshbach resonance effects to enhance the optical nonlinearity of two-dimensional polaritons in planar microcavities: by adjusting the cavity length it is in fact possible to bring the energy...
of a two-polariton state in resonance with a biexciton state.

Our work follows up on this pioneering proposal: we develop a simple analytical model of the biexciton Feshbach resonance and we analyze the nonlinear and quantum optical phenomenology of specific semiconductor devices. The theoretical model is inspired by the atom-molecule approach to strongly interacting degenerate quantum gases [10,11] and provides analytical formulas for both the energy-dependent scattering amplitude in a planar geometry and the single-polariton nonlinearity in confined geometries.

The main result of this letter is the characterization of the non-classical antibunching properties of the light emission from a micron-sized polariton dot [12–15]: the enhancement of the optical nonlinearity on the biexciton Feshbach resonance is anticipated to provide an efficient Feshbach blockade effect where the presence of a single polariton is able to block the entrance of a second one. As polariton dots do not involve quantum confinement of carriers, they are not subject to the fluctuations in the confinement length that are responsible for the typical carriers, they are not subject to the fluctuations in the confinement length over which the dynamicsof the exciton and cavity photons is described by a Hamiltonian term of the form [8]

\[
H_{BXC} = \frac{1}{2} \sum_{\sigma=\{\uparrow, \downarrow\}} \int d^2 R d^2 r \left\{ \sum_{i=(C,X)} \sum_{\sigma^\prime} \hat{\Psi}_{i,\sigma}^\dagger(r) \right. \\
\times \left[ \delta_{\sigma,\sigma^\prime} \left( E_i^o - \frac{\hbar^2 \nabla^2}{2 M_i} \right) + V_{i,\sigma,\sigma^\prime} \right] \hat{\Psi}_{i,\sigma^\prime}(r) \\
+ \hbar \Omega_R \sum_{\sigma} \left[ \hat{\Psi}_{C,\sigma}^\dagger(r) \hat{\Psi}_{X,\sigma}(r) + \hat{\Psi}_{X,\sigma}^\dagger(r) \hat{\Psi}_{C,\sigma}(r) \right] \\
\left. + \hat{\Psi}_B^\dagger(r) \left( E_B^o - \frac{\hbar^2 \nabla^2}{2 M_B} \right) \hat{\Psi}_B(r) \right\},
\]  

where \( E_{C,X,B}^o \) are the rest energies of a cavity photon, an exciton and a biexciton, respectively. \( M_{C,X,B} \) are the corresponding effective masses within a parabolic dispersion approximation. In the typical experimental systems, the exciton and biexciton masses \( M_{X,B} \) are comparable to the free electron mass (\( m_0 \)) and are much larger than that of the cavity photon (\( M_C \sim 10^{-5} m_0 \)). The spin variable \( \sigma = \uparrow, \downarrow \) identifies the polarization states of the cavity photon and of the exciton. The biexciton state in contrast is a singlet. The dark exciton states are not included in the model since they do not contribute to the physics we consider. \( V_{(C,X),\sigma^\prime}(r) \) denotes the — possibly spin-dependent [16]— confinement potential acting on the cavity photon and the exciton. \( \Omega_R \) is the Rabi frequency corresponding to the exciton-photon dipole coupling. A typical example of the momentum-space dispersion of the coupled exciton and cavity photons in a spatially homogeneous geometry \( V_{X,C} = 0 \) is shown in fig. 1(a): the exciton and cavity photon (dotted lines) are merged into the lower- (LP) and upper-polariton (UP) states (solid lines). On the scale of the figure, the dispersion of excitons is almost flat (dashed line).

A simple and transparent description of the physical process leading to the formation of a biexciton can be developed in terms of the giant oscillator strength model. Within this model, absorption of a cavity photon by an existing exciton leads to the creation of a biexciton via a Hamiltonian term of the form [8]

\[
H_{BXC} = \frac{1}{2} \sum_{\sigma=\{\uparrow, \downarrow\}} \int d^2 R d^2 r \left[ g(r) \hat{\Psi}_{B}^\dagger(R) \hat{\Psi}_{C,X}(R-r/2) + h.c. \right].
\]  

Even without a detailed knowledge of the microscopic structure of the biexciton, general physical arguments can be invoked to obtain important information on the coupling amplitude \( g(r) \): the optical process leading to the formation of a biexciton is limited to a spatial region around the existing exciton with a characteristic size on the order of the spatial size \( a_B \) of the biexciton. Its value scales as the amplitude of the wave function describing the relative motion of two excitons forming a biexciton \( a_B^{-1} h(r/a_B) \); here, \( h(s) \) is determined by the attractive exciton-exciton interaction potential and has a peak value and a spatial extension of order unity. The local value of the exciton creation matrix element is \( \Omega_R \). These two general observations lead us to the following generic form for the exciton-photon-biexciton coupling:

\[
g(r) = \frac{\hbar \Omega_R}{a_B} h(r/a_B).
\]  

If the exciton-photon Rabi frequency \( \Omega_R \) is comparable to or larger than the biexciton binding energy, we can tune the cavity in a way to ensure that the upper-polariton branch is far enough away in energy to be safely neglected and, at the same time, that the biexciton resonance is comparable to twice the energy of polariton states near the bottom of the lower-polariton dispersion (fig. 1(a), (c)). We can then rewrite the exciton-photon-biexciton coupling term \( H_{BXC} \) in terms of the
lower-polariton field only:

\[
H_{BXC} = \frac{1}{2} \sum_{\sigma, \sigma'} \int d^2 \mathbf{r} d^2 \mathbf{r}' \frac{\hbar \Omega_{\sigma, \sigma'}}{a_B} u_{LP}^X u_{LP}^C \bigg[ h(r/a_B) \hat{\Psi}^\dagger_{LP, \sigma}(\mathbf{r}) \hat{\Psi}_{LP, \sigma'}(\mathbf{r} - \mathbf{r}/2) \\
\times \hat{\Psi}_{LP, -\sigma}(\mathbf{r} + \mathbf{r}/2) + \text{h.c.} \bigg].
\]

Here, \( u_{LP}^X \) are the Hopfield coefficients quantifying the exciton and cavity photon component of the lower polariton; as we are focussing our interest to the lowest part of the lower-polariton dispersion only, the dependence of \( u_{LP}^X(k) \) on momentum \( k \) is neglected and their value at \( k = 0 \) is considered. The form (4) of the polariton-biexciton coupling emphasizes the crucial fact that the final product of the biexciton disintegration consists of a pair of lower polaritons rather than an exciton and a photon [17].

In general, there are additional exciton scattering processes that are independent of the biexciton resonance, analogous to the so-called background scattering length in the atomic case [5]. As is the case in the theory of cold quantum gases, the effective potential describing such scattering processes does not need to be the physical exciton-exciton potential. Any model potential can be chosen, provided it correctly reproduces the scattering physics in the region of parameters of interest once all direct and exchange processes are taken into account [18]. For the sake of simplicity, in the following we shall choose a repulsive, short-ranged form for the exciton-exciton potential \( v_{\sigma, \sigma'}(\mathbf{r} - \mathbf{r}') = v_{\sigma, \sigma'} \delta^2(\mathbf{r} - \mathbf{r}') \). If the relevant physics is limited to the lower-polariton branch, the same reasoning underlying (4) leads to a local polariton-polariton potential of rescaled intensity \( v_{\sigma, \sigma'} | u_{LP}^X |^4 \).

**Polariton-polariton scattering amplitude.** — Since our principal focus is on the biexciton-Feshbach resonance enhancement of the polariton-polariton interactions, we first neglect the contribution from the direct exciton-exciton scattering. Straightforward manipulations can be used to resum the diagrams contributing to the polariton-polariton scattering amplitude stemming from eq. (2).
Approximating the polariton dispersion with a parabola of mass $M_{LP}$, the scattering $T$-matrix element in the singlet channel has the resonant form

$$T_{11}(E) = \frac{|\bar{g}|^2}{E - E_{UV} + \frac{M_{LP}}{2 \hbar^2} |\bar{g}|^2 \log|E_{UV}/E| + i \frac{M_{LP}}{2 \hbar^2} |\bar{g}|^2}. \tag{5}$$

This expression for the $T$-matrix element was first discussed in [19] in the context of atom-atom scattering in two-dimensional geometries. The coefficient $\bar{g}$ is here defined as $\bar{g} = h \Omega_{R} a_B u_{G}^{*} u_{LP} \hbar$ with $\hbar = a_B^{-2} \int d^2 r \ h(\rho/a_B)$ is of order unity. No scattering occurs instead in the triplet channel.

From equation (5), it is easy to recognize several features typical of a Feshbach resonance in two-dimensional scattering. The radiative correction to the resonance energy has a logarithmic dependence on the energy of the form $\Delta_{rad} = M_{LP} |\bar{g}|^2 / (4 \hbar^2) \log|E_{UV}/E|$; the high-energy contribution to $\Delta_{rad}$ depends very much on the details of the system and has been absorbed into the ultraviolet parameter $E_{UV}$. The imaginary part of the denominator accounts for the radiative decay of the biexciton into a pair of propagating polaritons with an energy-independent linewidth $\Gamma_{rad} = M_{LP} |\bar{g}|^2 / 2 \hbar^2$ comparable to the radiative shift $\Delta_{rad}$. Inserting realistic parameters from GaAs-based structures, namely $h \Omega_{R} \approx 5 \text{ meV}$ and $a_B \approx 20 \text{ nm}$, the resulting radiative linewidth $\Gamma_{rad}$ into a pair of polaritons is on the order of a few $\mu$eV, somehow smaller than the value quoted in [20]. Additional contributions to the biexciton decay rate could be included in an effective $\Gamma_{tot} = \Gamma_{rad} + \Gamma_{add}$. A significant contribution to the additional decay rate $\Gamma_{add}$ may stem from dissipative coupling of the biexciton to bulk photons and/or interface polaritons. For simple planar geometries, we can estimate this contribution to be on the order of $10 \mu$eV. Other contributions such as Auger decay [21,22] can be reasonably expected to be much smaller.

At the mean-field level the dynamics of the coupled cavity photon, exciton and biexciton quantum fields can be described in terms of classical fields $\Psi_{C,X,B} = (\Psi_{C,X,B}(r))$. This approach has been already proven to be useful in studies of biexciton optics [8,23] as well as to describe the dynamics of atom-molecule conversion in ultracold atomic gases [10,11]. In contrast, we shall now focus on a spatially confined system where quantum fluctuations are instead strong and the mean-field approximation fails. Our goal is to demonstrate that the strong polariton-polariton interactions originating from the biexciton Feshbach resonance can translate into a significant polariton blockade effect [12] with peculiar antibunching features in the transmitted light.

**Zero-dimensional polariton dots.** - Zero-dimensional polariton dots [13–15] are characterized by a strong spatial confinement of the polariton wave function in all three dimensions, which gives rise to a series of discrete polariton states. In typical experimental set-ups, the in-plane confinement of polaritons is obtained by suitably designing a photonic structure whose discrete photonic eigenmodes have a small transverse mode profile. Confinement of the bare exciton or the biexciton is instead not necessary as it gets inherited from the photonic one.

If the confined polariton modes are well separated in energy as compared to their linewidth and all nonlinear couplings, we can restrict our attention to a single polariton mode. This assumption requires that we limit ourselves to geometries where the photon is confined to a length scale of the order of a $\mu$m. A moderate spatial anisotropy of the confinement is enough to have a sufficient spectral separation ($\gtrsim 1 \text{ meV}$) of the two lowest linearly polarized polariton modes [16], so that the spinor part of the (spinorial) wave function $\phi_{LP}(r, \sigma)$ becomes trivial $\phi_{LP}(r, \uparrow) = \phi_{LP}(r, \downarrow) = 1 / \sqrt{2} \phi_{LP}(r)$.

The creation operator of the single confined linearly polarized mode will be denoted by $\hat{b}$. Nonlinear processes occur via resonant scattering onto the intermediate biexciton state. Given the form (4) of the polariton-biexciton coupling and the very large value of the biexciton mass as compared to the polariton one, a single biexciton state of spatial wave function $\phi_{B}(r)$ proportional to the square $\phi_{LP}(r)^2$ will be involved in the dynamics. The corresponding biexciton creation operator will be denoted as $\hat{b}_c$. We emphasize that the biexciton states are not confined in our model; the wave function $\phi_{B}(r)$ is determined by the collective coupling amplitude to the relevant biexciton eigenmodes contributing to confined polariton-polariton scattering.

Limiting our attention to these polariton and biexciton states, the resulting Hamiltonian can be written in the concise form

$$H = E_{LP} \hat{b}_c \hat{b} + E_B \hat{b}_c \hat{b} + \frac{1}{2} V \hat{p}_c \hat{p}_c \hat{b}_c \hat{b}_c + G \left[ \hat{b}_c \hat{p}_c + \hat{b}_c \hat{p}_c \hat{b}_c \right]. \tag{6}$$

Here, $E_{LP}$ and $E_B$ are the energies of the confined lower-polariton and biexciton states, respectively. The nonlinear interaction coefficients $V$ and $G$ are evaluated by inserting the lower-polariton $\phi_{LP}(r, \sigma)$ and biexciton $\phi_{B}(r)$ wave functions into the interaction Hamiltonian and neglecting contributions from all other states.

$$V = \frac{1}{4} \sum_{\sigma, \sigma'} |u_{LP}^{X}(r)|^4 \int d^2 r |\phi_{LP}(r)|^4, \tag{7}$$

$$G = \hbar \Omega_{R} \frac{a_B}{u_{LP}^{X}} \int d^2 r d^2 R \ h(r/a_B) \phi_{B}(R) \times \phi_{LP}(R - r/2, \uparrow) \phi_{LP}(R + r/2, \downarrow) \approx \frac{1}{2} \hbar \Omega_{R} a_B u_{LP}^{X} \int d^2 r |\phi_{LP}(r)|^4 \ h \left[ \frac{1}{2} |\phi_{LP}(r)|^4 \right]^{1/2}. \tag{8}$$

An important direct consequence of (7) and (8) is that the background interaction energy decreases as $1 / \ell^2$.

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for increasing confinement size $\ell$, while the polariton-biexciton coupling has a slower scaling proportional to $1/\ell$ only.

Previous calculations of the background exciton-exciton scattering [18] suggest that the background interaction energy $v_{bg,\sigma'}$ is of the order of $R^*a_X^2$, where $R^*$ is the exciton Rydberg energy (on the order of 10 meV for standard GaAs quantum wells) and $a_X$ is the exciton Bohr radius (on the order of 10 nm). As the radius $a_B \approx 20$ nm of Wannier bieexcitons is of the same order as the exciton radius $a_X$ and the Rabi energy $\Omega_R \approx 2$ meV is typically of the same order as the exciton Rydberg $R^*$, the resonant contribution $G$ is expected to dominate over the background scattering contribution $V$ as soon as the polariton dot is much larger than the physical size of excitons and bieexcitons $\ell \gg a_X, a_B$. As the confinement length $\ell$ of typical polariton dots is on the order of 1 $\mu$m, this latter condition is generally very well satisfied.

Inserting these values of the parameters in (8), one obtains a value of $G$ on the order of 30 $\mu$eV, which is extremely promising in view of applications: confined polariton states with a linewidth of 70 $\mu$eV have been recently demonstrated [15]. On the other hand, this value of the nonlinear coupling is safely below the energy distance from the neighboring polariton states in the dot; the typical energy splitting between confined polariton states in a box potential with a transverse size of the order of 1 $\mu$m are in fact on the order of a few meV, which guarantees a posteriori the validity of our single-mode calculations. As compared to the original proposal [12], the use of the biexciton Feshbach resonance allows us to relax the constraint on the size of the polariton dot.

**Few-body states.** – As the Hamiltonian (6) conserves the total number of excitation $N_{\text{tot}} = N_{\text{LP}} + 2N_B$, it can be easily diagonalized within each $N_{\text{tot}}$ manifold, giving a ladder of mixed levels. An example of this level ladder is shown in fig. 1(b) for the most remarkable case, $2E_{\text{LP}} + V = E_B$, where the biexciton is resonance with the (slightly shifted) polariton states. The $N_{\text{tot}} = 0$ manifold consists of the vacuum state $|\text{vac}\rangle$ only, while the $N_{\text{tot}} = 1$ manifold consists of a single one-polariton state $|p\rangle$ with energy $E_{\text{LP}}$. The $N_{\text{tot}} = 2$ manifold contains a doublet of eigenstates which are the symmetric and antisymmetric combinations of a bieexciton $|B\rangle$ and a pair of polaritons with identical linear polarization, $|pp\rangle$. The energies of these states are symmetrically split by an amount $\pm \sqrt{2G}$, typically much larger than $V$. In this way, the harmonicity of the level ladder is broken: the transitions from the $N_{\text{tot}} = 2$ to the $N_{\text{tot}} = 1$ manifold occur in fact at energies $E_{\text{LP}} \pm \sqrt{2G}$ different from the energy $E_{\text{LP}}$ of the $N_{\text{tot}} = 1$ to $N_{\text{tot}} = 0$ transition.

These simple considerations straightforwardly extend to the eigenstates forming the higher $N_{\text{tot}} > 2$ manifolds, which then consist of linear superpositions of states with $N_{\text{tot}} = 2n$ polaritons and $n$ bieexcitons ($0 \leq n \leq N_{\text{tot}}/2$) with suitable weights. In particular, it is interesting to note that, differently from the atomic case, the absence of deep molecular states other than the biexciton state should suppress all those three- and many-body recombination processes that generally limit the lifetime of strongly interacting atomic Bose gases [24].

Far from resonance $|2E_{\text{LP}} - E_B| \gg G$, the polariton and bieexciton states are no longer efficiently mixed by the polariton-biexciton coupling and their energies approach their uncoupled values. In this regime, the harmonicity of the level ladder is only broken by a much weaker background scattering term of magnitude $V \ll G$. A complete picture of the anticrossing between the bieexciton and polariton states in the $N_{\text{tot}} = 2$ manifold is shown in fig. 1(d): while in the atomic case the difference in magnetic moment of the molecular and atomic states allows to tune their relative energy by means of an external magnetic field, here the resonance between the polaritons and the biexciton can be obtained by suitably adjusting the cavity length. As it is shown in fig. 1(c), (d), the polariton energy strongly depends on the cavity photon energy (and in turn on the cavity length), while the biexciton energy only depends on the microscopic structure of the quantum well.

**Optical signatures of polariton blockade.** – The pronounced anharmonicity of the level scheme of fig. 1(b) has direct implications for the optical properties of the polariton dot. For the sake of simplicity we restrict here our attention to the $G > \Gamma, \Gamma_{\text{add}}$ limit, where $\Gamma$ is the radiative linewidth of the polariton states and $\Gamma_{\text{add}}$ is the additional, non-radiative contribution to the biexciton linewidth. The driving laser is assumed to be fully coherent, weak and resonant with the transition to the one polariton state $|p\rangle$, which then gets effectively populated. On the other hand, the Feshbach-resonance-induced nonlinearity shifts both states in the $N_{\text{tot}} = 2$ manifold far from resonance with the laser by an amount $\sqrt{2G}$ larger than the linewidth of the relevant states. As a consequence of this polariton blockade effect, the $N_{\text{tot}} = 2$ manifold remains almost unpopulated: the first polariton is able to effectively block the entrance of the second. In a transmission geometry where only light emitted by the polariton dot is collected, this polariton blockade effect is clearly visible as a complete anticorrelation between transmitted photons\(^{1}\). The efficiency of this effect is shown in fig. 1(e) (solid line), where the second-order correlation function of the transmitted light is plotted as a function of the exciton-biexciton coupling $G$: as soon as $G$ is larger than the damping rates $\Gamma, \Gamma_{\text{add}}$, the emitted light is strongly antibunched.

A different strategy to obtain strongly antibunched light involves the use of dissipative nonlinearities instead of

\(^{1}\)If the polarization degeneracy is not lifted, the Feshbach enhancement of the polariton-polariton interactions is restricted to the singlet channel, which limits the blockade effect to polaritons of opposite polarizations. This dramatically suppresses the antibunching features, but leaves interesting polarization correlations in the emitted light.
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simple way to obtain a strongly antibunched emission also

In our system, strong dissipative nonlinearities provide a

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conservative ones as done so far. A recent experiment with

atomic gases [25] has demonstrated how dissipative collis-

sional processes offer a viable alternative route to create

strongly correlated states. The possibility of extending this

idea to the photonic case was recently discussed in [26].

In our system, strong dissipative nonlinearities provide a

simple way to obtain a strongly antibunched emission also

in the Γ_{add} \gg G \gg Γ regime: in this case, the energy spac-

ing between the levels remains almost constant, but the

harmonicity of the level ladder is broken by the strongly

varying dissipation rate. In particular, for sufficiently large

values of exciton-biexciton coupling G, the decay rate of

the two-polariton \(|pp\rangle\) state \(Γ_2 = 2Γ + 2G^2/Γ_{add}\) can be

much larger than the decay rate Γ of the one-polariton \(|p\rangle\)

state. As a result, the excitation of the system under a

weak and coherent pump is again limited to the

\(N_{tot} = 1\) manifold, which produces a stream of antibunched photons

in the emission. The efficiency of this scheme to obtain

antibunched light is shown in fig. 1(e) (dashed line): in this case, antibunching sets in as soon as \(G \gg √Γ_{add}Γ\).

Which one of these schemes is the most suitable one

depends on the properties of the specific system under

consideration. The decay rate Γ of the lowest-energy cavity polaritons is determined exclusively by the cavity

losses. While the lowest value observed to date is \(Γ \sim 70 \mu eV\), there are no physical constraints that prohibit

reaching much smaller values by increasing the cavity

quality factor. On the other hand, the intrinsic biexciton

decay rate \(Γ_{add}\) into bulk photons or surface polaritons can be estimated to be on the order of \(10 \mu eV\). Unless

some specific effect is made to suppress it in a suitably

designed photonic structure, this biexciton decay rate is

most likely to set the ultimate limit to the antibunching

that is achievable in polariton dots.

Conclusions. – We have theoretically investigated the

physics of coupled exciton, photon and biexciton quantum

fields in planar semiconductor microcavities. The Fesh-

bach resonance on the biexciton state can be exploited to

dramatically enhance the amplitude of polariton scattering

processes. A strong photon antibunching is predicted for

the emitted light from strongly confined polariton dots.

If our predictions are verified experimentally, biexciton

physics will become the cornerstone for a new generation of

opto-electronic devices working at the single-photon

level, with a number of applications from fundamental

science to quantum information processing.

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