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MISCELLANEOUS

On the Mechanism of the Maintenance of Rabi Oscillations in the System of Exciton Polaritons in a Microcavity

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A physical mechanism of the implementation of undamped Rabi oscillations in the system of exciton polaritons in a semiconductor microcavity in the presence of nonresonant pumping has been proposed. Various mechanisms of the stimulated scattering of excitons from the reservoir have been considered. It has been shown that undamped oscillations of the population of the photon component of the condensate can be caused by the feeding of a coherent Rabi oscillator owing to the pair scattering of excitons from the reservoir to the ground state. The effect should be observed in spite of a more intense relaxation of polaritons of the upper branch observed experimentally.

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INTRODUCTION

The interaction of the photon field of a semiconductor microcavity with excitons localized in a quantum well in the strong coupling regime is responsible to the formation of two new eigenmodes—polaritons of the upper and lower branches. Owing to their quantum nature, exciton polaritons can macroscopically occupy the ground energy state and form a polariton condensate [1], which is characterized by a high degree of temporal and spatial coherence.

One of the manifestations of the coherent properties of exciton polaritons is the existence of Rabi oscillations, i.e., the periodic exchange of populations between the exciton and photon subsystems of the condensate. The frequency of oscillations is determined by the energy splitting between polaritons of the upper and lower branches. Processes occurring "inside" the polariton condensate, i.e., between its exciton and photon subsystems, in particular, Rabi oscillations, are currently under active investigation in view of their possible applications in integrated optical circuits [2–5].

Rabi oscillations can be detected by the pump–probe method, which allows measurements of beatings of an optical signal from a microcavity [6, 7] or oscillations of the magnetic moment of the exciton component [8]. Owing to relaxation processes appearing because of the open dissipative nature of the polariton condensate, Rabi oscillations in real structures are observed in a very short time (about several picoseconds [8]). A mechanism for an increase in the damping time of Rabi oscillations owing to the stimulated scattering to the Rabi oscillator from the polariton reservoir formed by incoherent pumping was proposed in [2]. This effect was then experimentally confirmed in [6]. However, an important problem of the possibility of implementing self-maintained permanent Rabi oscillations has not yet been solved in the cited works. This problem is studied in this work.

POLARITON CONDENSATE IN THE PRESENCE OF NONRESONANT PUMPING

We consider a spatially uniform condensate where the component of the wave vector in the plane of the microcavity is zero, $\mathbf{k} = 0$. For simplicity, the spin degrees of freedom are neglected. We describe the dynamics of the condensate in the mean field approximation by the amplitudes of the photon (ϕ) and exciton (χ) components, which depend only on the time and satisfy the nonlinear Schrödinger equation. In the presence of nonresonant pumping, the condensate interacts with the incoherent reservoir, whose dynamics satisfies the Boltzmann equation [9]:

$$\dot{\chi} = \frac{1}{2}(p_{\rm ex} - \gamma_{\rm ex})\chi + i\delta\chi - i\Omega\phi - ig_{\rm BS}\chi, \quad (1a)$$

$$\dot{\phi} = -\frac{1}{2}\gamma_{\rm ph}\phi - i\Omega\chi - iF_{\rm p}(t), \quad (1b)$$

$$\dot{N} = P - \gamma_{\rm R}N - p_{\rm ex}|\chi|^2. \quad (1c)$$

Here, 2Ω is the Rabi splitting between the exciton ($\omega_{\rm ex}$) and photon ($\omega_{\rm ph}$) modes of the microcavity at zero detuning $\delta = \omega_{\rm ph} - \omega_{\rm ex}$; $\gamma_{\rm ex}$ and $\gamma_{\rm ph}$ describe the damping of the exciton and photon modes, respectively; the quantity $g_{\rm BS} = g_{\rm c}|\chi|^2 + g_{\rm R}N$ describes the elastic scattering of excitons in the ground state by each other (with the interaction constant $g_{\rm c}$) and by excitons of the reservoir ($g_{\rm R}$). Both of these processes result in a blue shift of the energy level of excitons. Rabi oscillations are excited by a femtosecond optical pulse $F_{\rm p}(t)$, whose duration is much smaller than the period of oscillations $\sim\Omega^{-1}$. The action of such a pulse can be described by choosing the initial conditions in the form $\chi_0 = 0$ and $\phi_0 = \sqrt{\int F(t)dt}$.

The reservoir of incoherent excitons characterized by the density N is formed by an external spatially uniform cw-pump P. The last term in Eq. (1c) describes the depletion of the reservoir owing to stimulated transitions of excitons to the Rabi oscillator specified by Eqs. (1a) and (1b). The effect of all other processes responsible for the departure of excitons from the reservoir without the filling of the ground state is described by the generalized term $\gamma_{\rm R}N$.

We consider two mechanisms of pumping of the exciton mode from the reservoir, specifying the rate of this process in the form

$$p_{\rm ex} = R_{\rm phon}N + R_{\rm ex-ex}N^2|\chi|^2, \quad (2)$$

where the first term describes the process of energy relaxation of the exciton from the reservoir to the ground state owing to scattering by an acoustic phonon [9]. The rate of such transitions is proportional to the density of particles in the reservoir with the coefficient $R_{\rm phon}$. The second considered mechanism is the pair scattering of excitons with the momenta $\hbar\mathbf{k}$ and $-\hbar\mathbf{k}$ to the state with $\mathbf{k} = 0$. Since such a process involves two incoherent excitons, its rate is proportional to N^2 and to the number of excitons $|\chi|^2$ in the ground state. In order to estimate the coefficient $R_{\rm ex-ex}$, it seems reasonable to assume that this process, as well as any nonlinear process, is less probable than scattering by phonons at a weak pumping and plays a noticeable role only at a quite strong pumping P.

In this work, plots are drawn with the following parameters of the system (see [3]), which are characteristic of GaAs-based microcavities: $\gamma_{\rm ex} = 0.01$ ps^{-1}, $\gamma_{\rm ph} = 0.1$ ps^{-1}, $\hbar\Omega = 2.5$ meV, $\gamma_{\rm R} = 0.003$ ps^{-1}, $g_{\rm c} = 0.009$ μm^2 ps^{-1}, $g_{\rm R} = 2g_{\rm c}$, $\hbar R_1 = 0.01$ μm^2 meV, and $\hbar R_2 = 0.001$ μm^6 meV.

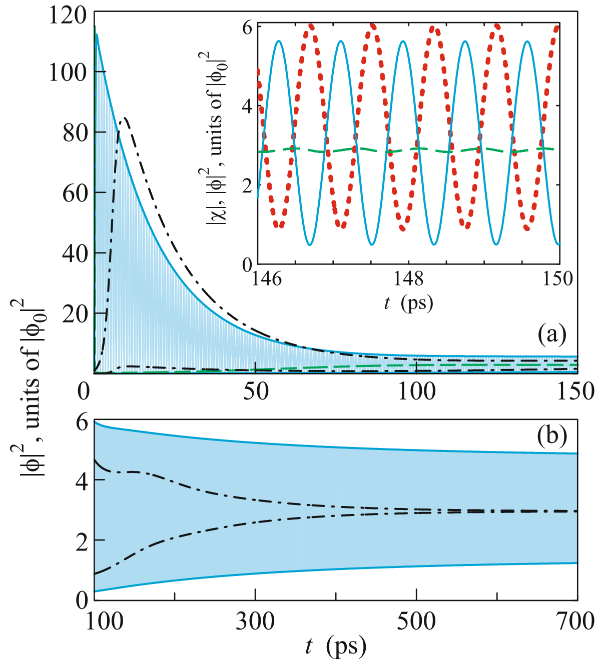

Fig. 1. (Color online) (a) Time dependence of the population of the exciton mode $|\phi|^2$ normalized to the initial value $|\phi_0|^2$ for $\delta = 0.2\Omega$ and $P = 0.25$ μm^{-2} ps^{-1}. The dashed line corresponds to the population of the reservoir N. The dash-dotted line is the envelope of oscillations $|\phi|^2$ in the limit $R_{\rm ex-ex} = 0$. The inset shows the modes (solid line) $|\chi|^2$ and (dashed line) $|\phi|^2$ in a short time interval. (b) Same as in panel (a) but beginning with the time $t = 100$ ps.

Below, we analyze the conditions of establishing undamped Rabi oscillations, which corresponds to the solution of Eqs. (1a) and (1b) in the form

$$\chi = \chi_1\exp(i\omega_1 t) + \chi_2\exp(i\omega_2 t), \quad (3a)$$

$$\phi = \phi_1\exp(i\omega_1 t) + \phi_2\exp(i\omega_2 t). \quad (3b)$$

When such a solution exists, the populations of the exciton and photon components undergo Rabi oscillations at the frequency $\omega_1 - \omega_2$ (see inset in Fig. 1a), where ω_1 and ω_2 are the natural (polariton) frequencies of the system, which are measured from the frequency of the photon mode of the microcavity $\omega_{\rm ph}$. Then, the amplitudes $\chi_{1,2}$ and $\phi_{1,2}$ are proportional to the fractions of the upper (subscript 2) and lower (subscript 1) polariton states in the exciton and photon modes, respectively.

In the limit $\Omega \gg p_{\rm ex}$, which is valid at strong coupling, processes of the redistribution of the populations between the components of the Rabi oscillator (exciton and photon modes) are much faster than processes of pumping of the exciton mode from the reservoir. As a result, we can adiabatically exclude Eq. (1c), taking N as constant at the time scale of the period of Rabi oscillations ~Ω^{-1}. Indeed, as is shown in Fig. 1a, the population of the reservoir (dashed line) in the regime of undamped oscillations oscillates with a small amplitude around a certain average value $\bar{N}$. This approximation is valid at relatively small pumps, when $p_{ex}|\chi|^2 \ll \gamma_R N$ (see Eq. (1c)). In this work, we use this approximation, considering the system not too far from the threshold of the formation of the condensate. The depletion of the reservoir in the regime of strong pump can result in undesirable damping of Rabi oscillations.

Since the system is dissipative, the eigenfrequencies $\omega_{1,2}$ are generally complex, corresponding to damping or amplification of Rabi oscillations. We seek solutions for which $\omega_{1,2}$ are real and $\chi_{1,2}$ and $\phi_{1,2}$ are nonzero. In this case, the conditions of the achievement of undamped oscillations are determined by a particular mechanism of scattering from the reservoir to the main mode. We consider the effect of two described mechanisms separately neglecting the blue shift for simplicity.

PUMPING OF THE EXCITON MODE OWING TO THE SCATTERING OF EXCITONS FROM THE RESERVOIR ON PHONONS

We substitute solution (3) into Eqs. (1a) and (1b) at $p_{ex} = R_{phon}\bar{N}$. Then, we separate the real terms from imaginary ones in the equations under the assumption that $\omega_{1,2}$ are real. As a result, we obtain

$$\chi_{1,2}[R_{phon}\bar{N}] = \chi_{1,2}\left(\gamma_{ex} + \frac{\Omega^2\gamma_{ph}}{\gamma_{ph}^2/4 + \omega_{1,2}^2}\right), \quad (4a)$$

$$\omega_{1,2}^3 - \delta\omega_{1,2}^2 - \omega_{1,2}(\Omega^2 - \gamma_{ph}^2/4) - \delta\gamma_{ph}^2/4 = 0. \quad (4b)$$

Relation (4a) is the condition of balance between pump and loss for polaritons in the upper and lower branches. Balance can be achieved for both branches simultaneously only at $\omega_1 = -\omega_2$, which is possible only if $\delta = 0$ (see Eq. (4b)). In planar microcavities studied experimentally, the dissipation rate of polaritons in the upper branch is usually much higher than that in the lower branch because an additional scattering channel exists in the upper branch.

At exciton–photon resonance ($\delta = 0$), excitons and photons make the same contribution to polariton states. In this case, losses in the photon mode for both branches are compensated by the pumping of the exciton component owing to scattering from the reservoir, as directly follows from Eqs. (4a) and (4b):

$$R_{phon}\bar{N} - \gamma_{ex} = \gamma_{ph}. \quad (5)$$

At $\delta \neq 0$, a more exciton-like state will be excited strongly.

We emphasize that the existence of undamped oscillations implies that the eigenfrequencies of the system of Eqs. (1a) and (1b) are real. When the blue shift is neglected and under condition (5), the system of Eqs. (1a) and (1b) can be written in the form

$$i\frac{d\phi}{dt} = \Omega\chi - i\frac{\gamma_{ph}}{2}\phi, \quad i\frac{d\chi}{dt} = \Omega\phi - i\frac{\gamma_{ph}}{2}\chi, \quad (6)$$

which is similar to the representation of the system of coupled oscillators (optical waveguides) in the $\mathcal{PT}$-symmetric systems [10, 11]. In such systems, because of the spatial asymmetry of the imaginary part of the refractive index of the medium, loss in one of the waveguides is exactly compensated by amplification in another waveguide. As a result, the optical pulse propagates in the medium without loss. In our case, the exciton and photon modes serve as waveguides. Thus, the existence of real eigenfrequencies in the system under consideration can be interpreted as the manifestation of properties of pseudo-Hermitian systems (see [12]) by the polariton condensate in the presence of the reservoir. This can be detected in the form of undamped Rabi oscillations.

PUMPING OF THE EXCITON MODE OWING TO EXCITON–EXCITON SCATTERING FROM THE RESERVOIR

When exciton–exciton scattering dominates in the system, $p_{ex} = R_{ex-ex}N^2|\chi|^2$, the condition of existence of solution (3) changes significantly. The substitution of Eq. (3) into Eqs. (1a) and (1b) gives an analog of the condition of balance given by Eq. (4a):

$$R_{ex\text{-}ex}\bar{N}^2(|\chi_{1,2}|^2 + 2|\chi_{2,1}|^2) = \gamma_{ex} + \frac{\Omega^2\gamma_{ph}}{\gamma_{ph}^2/4 + \omega_{1,2}^2}. \quad (7)$$

At the same time, Eq. (4b) does not change. Since Eq. (7) includes the average population of the reservoir $\bar{N}$, it should be treated as the equation of balance valid at times much larger than the period of Rabi oscillations. For this reason, terms oscillating at the Rabi frequency $\omega_1 - \omega_2$ can be neglected when deriving Eq. (7).

It is worth noting that the rate of pumping of the polariton state will depend not only on $\bar{N}$ but also on the populations of *both* polariton branches, in particular, on the degree of their contribution to the exciton modes $|\chi_1|^2$ and $|\chi_2|^2$ (see Eq. (7)). The recharge of the population of the coherent Rabi oscillator owing to exciton–exciton scattering from the reservoir generally increases the robustness of the system to imbalance of dissipation in the upper and lower polariton branches and makes it possible to reach permanent Rabi oscillations in a wide range of the parameters.

Figure 1 demonstrates the dynamics of establishing undamped oscillations for the population of the photon mode $|\phi|^2$ in the system, which can be measured experimentally. Because of the strong nonlinearity of the pump mechanism based on exciton–exciton scattering, the envelope of oscillations $|\phi|^2$ undergoes a sharp jump in the first several picoseconds, which is accompanied by the fast depletion of the reservoir (dashed line). Oscillations with a constant amplitude between the exciton and photon components are established in the system after approximately 100 ps (at the chosen parameters of the system) and occur in antiphase (see inset in Fig. 1a).

To plot Fig. 1, the reservoir particle generation rate P was chosen significantly above the threshold. For this reason, the pump mechanism based on the exciton–exciton scattering dominates in the system. The dash-dotted line in Fig. 1 corresponds to the limit $R_{ex-ex} = 0$. In this case, since $\delta \neq 0$, the population of the photon component at large times becomes stationary; i.e., Rabi oscillations damp (see Fig. 1b).

It is noteworthy that the mechanism of maintenance of Rabi oscillations owing to the exciton–exciton scattering is effective in a limited detuning range $|\delta| < \delta_c$. At a certain (e.g., positive) detuning $\delta = \delta_c$, imbalance between loss and pump for polariton states becomes so large that, to compensate it, the quantity $|\chi_2|^2$ should be equal to the total number of excitons in the system $|\chi|^2$. In this case, $\chi_1 = 0$ and Rabi oscillations disappear (see Eq. (3)).

To determine δ_c in the limit $p_{ex} = R_{ex-ex}N^2 |\chi|^2$, we neglect the term $\delta\gamma_{ph}^2/4$ in Eq. (4b) in the strong coupling regime. Then, solving this equation together with Eq. (7) under the condition $|\chi_1|^2 = 0$, we obtain

$$\delta_c \simeq \pm\left[\frac{\Omega^2}{4\gamma_{ph}}\left(\gamma_{ex} - 8\gamma_{ph} + 3\sqrt{\gamma_{ex}^2 + 8\gamma_{ph}^2}\right)\right]^{1/2}. \quad (8)$$

The condition $|\chi_2|^2 = 0$ provides a difference only in the sign of δ_c. At the chosen parameters, $\delta_c \simeq 0.38\Omega = 1.49\ \text{ps}^{-1}$.

Both of the above mechanisms of pumping of the exciton mode from the reservoir should be taken into account simultaneously in real systems. Combining Eqs. (4) and (7), we arrive at the generalized condition of balance

$$R_{ex-ex}\bar{N}^2(|\chi_{1,2}|^2 + 2|\chi_{2,1}|^2) + R_{phon}\bar{N} = \gamma_{ex} + \frac{\Omega^2\gamma_{ph}}{\gamma_{ph}^2/4 + \omega_{1,2}^2}. \quad (9)$$

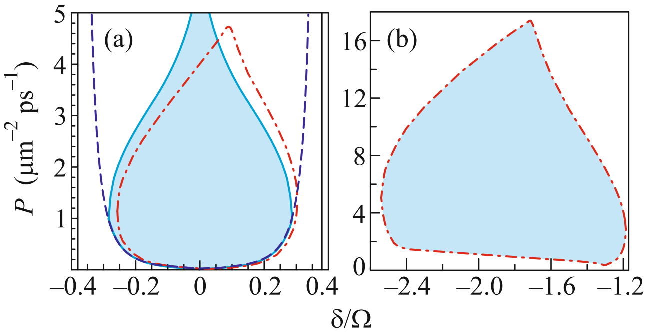

Fig. 2. (Color online) Diagram of the existence of undamped Rabi oscillations (shaded region) (a) without and (b) with allowance for the relaxation of the upper polariton branch at $\gamma' = 0.4\ \text{ps}^{-1}$. The dashed lines in panel (a) correspond to the value determined by Eq. (8). The dash-dotted line specifies the region of existence of oscillations, which is numerically determined with allowance for the blue shift of the energy level of the exciton.

The $\bar{N}$ value can be found from the condition (see Eqs. (1c) and (3))

$$P - \gamma_R\bar{N} - R_{phon}\bar{N}(|\chi_1|^2 + |\chi_2|^2) - R_{ex-ex}\bar{N}^2(|\chi_1|^4 + |\chi_2|^4 + 4|\chi_1|^2|\chi_2|^2) = 0. \quad (10)$$

To determine the region of the parameters at which the regime of permanent oscillations is achieved, we solve Eqs. (9) and (10) self-consistently setting successively $|\chi_1|^2 = 0$ and $|\chi_2|^2 = 0$. The dashed line in Fig. 2a corresponds to this solution. Immediately near the threshold P^{thr}, scattering by phonons dominates in the system. Consequently, undamped oscillations exist only near resonance. As the pump P increases, this region is expanded and tends to the limit predicted by Eq. (8) (vertical dashed lines). At quite large P values, Eq. (9), which is written for quantities averaged over the period of Rabi oscillations, is no longer valid. In this limit, the reservoir is depleted and the amplitude of its fluctuations induced by Rabi oscillations becomes comparable with the average value $\bar{N}$, and permanent Rabi oscillations disappear (cf. [13]). The solid curve in Fig. 1a limits the region of existence of undamped oscillations, which is determined from the numerical solution of system (1) (at $g_{BS} = 0$). It is seen in this figure that our analysis describes well the behavior of the system at not overly large P values. It is also worth noting that the Mott transition occurs at a very large pump, excitons are transformed to an electron–hole plasma, and all effects associated with them disappear.

INCLUSION OF THE EXCITON ENERGY BLUE SHIFT AND ADDITIONAL RELAXATION OF THE UPPER POLARITON BRANCH

We consider the problem of the maintenance of Rabi oscillations with the inclusion of the blue shift of the energy level of the exciton and additional relaxation of polaritons in the upper branch, which, as is known [14], plays an important role in the damping of Rabi oscillations in real structures.

The above consideration can be repeated with the inclusion of the term $g_{\rm BS}$, which is responsible for the blue shift. In this case, the structure of Eqs. (9) and (10) does not change, but the quantity δ should be replaced by the effective detuning

$$\tilde{\delta}_{1,2} = \delta - g_{\rm R}\bar{N} - g_{\rm c}(|\chi_{1,2}|^2 + 2|\chi_{2,1}|^2). \qquad (11)$$

Thus, weak interaction between excitons leads only to a change in the range of the parameters at which undamped oscillations appear (dash-dotted line in Fig. 2a) but does not affect the possibility of their detection.

To take into account the effect of additional relaxation of polaritons in the upper branch, we introduce polariton states in the mean field approximation:

$$a_{\rm LB} = C_x\chi - C_p\phi,\ a_{\rm UB} = C_x\phi + C_p\chi, \qquad (12)$$

where

$$C_{x,p} = \frac{1}{\sqrt{2}}\left(1 \pm \frac{\delta}{\sqrt{4\Omega^2 + \delta^2}}\right)^{1/2}$$

are the Hopfield coefficients and $|a_{\rm LB}|^2$ and $|a_{\rm UB}|^2$ are the number densities of polaritons in the lower and upper branches, respectively.

Using definition (12), we can represent Eqs. (1a) and (1b) in the polariton basis. Additional relaxation of polaritons in the upper branch at the rate γ' is taken into account phenomenologically by adding the term $-\gamma' a_{\rm UB}/2$ to the resulting equation for $\dot{a}_{\rm UB}$. Then, we perform the inverse transition to the exciton–photon basis and obtain the system

$$\dot{\chi} = \frac{1}{2}(p_{\rm ex} - \tilde{\gamma}_{\rm ex})\chi + i\delta\chi - i\tilde{\Omega}\phi - ig_{\rm BS}\chi, \qquad (13a)$$

$$\dot{\phi} = -\frac{1}{2}\tilde{\gamma}_{\rm ph}\phi - i\tilde{\Omega}\chi, \qquad (13b)$$

where $\tilde{\gamma}_{\rm ex} = \gamma_{\rm ex} + \gamma' C_p^2$, $\tilde{\gamma}_{\rm ph} = \gamma_{\rm ph} + \gamma' C_x^2$, and $\tilde{\Omega} = \Omega - i\gamma'/2C_xC_p$.

The numerical solution of system (13) makes it possible to find the region of existence of undamped oscillations with the inclusion of additional relaxation of polaritons in the upper branch and of the shift of the energy of excitons (see Fig. 2b). It is seen in Fig. 2b that undamped oscillations exist only at negative detunings quite large in absolute value. Polaritons in the upper branch in this limit become strongly exciton-like and acquire a large fraction of pump of the exciton state from the reservoir, which allows the compensation of large loss in the upper polariton branch.

CONCLUSIONS

It has been shown that undamped Rabi oscillations can be established in the exciton–photon system in a microcavity. This regime can be achieved at various exciton–photon detunings, as well as in the presence of additional relaxation in the upper polariton branch. The proposed mechanism of compensation of loss owing to pair exciton scattering from the reservoir is flexible and efficient and can promote the establishment of undamped Rabi oscillations in a wide region of parameters.

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REFERENCES

1. J. Kasprzak, M. Richard, S. Kundermann, A. Baas, P. Jeambrun, J. M. J. Keeling, F. M. Marchetti, M. H. Szymańska, R. Andre, J. L. Staehl, V. Savona, P. B. Littlewood, B. Deveaud, and L. S. Dang, Nature (London) **443**, 409 (2006).
2. S. S. Demirchyan, I. Yu. Chestnov, A. P. Alodjants, M. M. Glazov, and A. V. Kavokin, Phys. Rev. Lett. **112**, 196403 (2014).
3. T. C. H. Liew, Y. G. Rubo, and A. V. Kavokin, Phys. Rev. B **90**, 245309 (2014).
4. N. S. Voronova, A. A. Elistratov, and Yu. E. Lozovik, Phys. Rev. Lett. **115**, 186402 (2015).
5. A. V. Kavokin, A. S. Sheremet, I. A. Shelykh, P. G. Lagoudakis, and Y. G. Rubo, Sci. Rep. **5**, 12020 (2015).
6. L. Dominici, D. Colas, S. Donati, J. P. Restrepo Cuartas, M. de Giorgi, D. Ballarini, G. Guirales, J. C. Lopez Carreno, A. Bramati, G. Gigli, E. del Valle, F. P. Laussy, and D. Sanvitto, Phys. Rev. Lett. **113**, 226401 (2014).

7. D. Colas, L. Dominici, S. Donati, A. A. Pervishko, T. C. H. Liew, I. A. Shelykh, D. Ballarini, M. de Giorgi, A. Bramati, G. Gigli, E. del Valle, F. P. Laussy, A. V. Kavokin, and D. Sanvitto, Light: Sci. Appl. **4**, e350 (2015).
8. A. Brunetti, M. Vladimirova, D. Scalbert, M. Nawrocki, A. V. Kavokin, I. Shelykh, and J. Bloch, Phys. Rev. B **74**, 241101(R) (2006).
9. M. Wouters and I. Carusotto, Phys. Rev. Lett. **99**, 140402 (2007).
10. A. A. Zyablovskii, A. P. Vinogradov, A. A. Pukhov, A. V. Dorofeenko, and A. A. Lisyanskii, Phys. Usp. **57**, 1063 (2014).
11. H. Ramezani, T. Kottos, R. El-Ganainy, and D. N. Christodoulides, Phys. Rev. A **82**, 043803 (2010).
12. J. Y. Lien, Y.-N. Chen, N. Ishida, H.-B. Chen, C.-C. Hwang, and F. Nori, Phys. Rev. B **91**, 024511 (2015).
13. C. Sturm, D. Solnyshkov, O. Krebs, A. Lemaitre, I. Sagnes, E. Galopin, A. Amo, G. Malpuech, and J. Bloch, Phys. Rev. B **91**, 155130 (2015).
14. B. Sermage, S. Long, I. Abram, J. Y. Marzin, J. Bloch, R. Planel, and V. Thierry-Mieg, Phys. Rev. B **53**, 16516 (1996).

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