Backscattering Suppression in Supersonic 1D Polariton Condensates

D. Tanese,[1] D. D. Solnyshkov,[2] A. Amo,[1] L. Ferrier,[1] E. Bernet-Rollande,[1] E. Wertz,[1] I. Sagnes,[1] A. Lemaître,[1] P. Senellart,[1] G. Malpuech,[2] and J. Bloch[1]

[1]*Laboratoire de Photonique et Nanostructures, LPN/CNRS, Route de Nozay, 91460 Marcoussis, France*
[2]*Institut Pascal, PHOTON-N2, Clermont Université, University Blaise Pascal, CNRS, 24 avenue des Landais, 63177 Aubière cedex, France*

(Received 26 August 2011; published 19 January 2012)

We investigate the effect of disorder on the propagation of one-dimensional polariton condensates in semiconductor microcavities. We observe a strong suppression of the backscattering produced by the imperfections of the structure when increasing the condensate density. This suppression occurs in the supersonic regime and is simultaneous to the onset of parametric instabilities which enable the "hopping" of the condensate through the disorder. Our results evidence a new mechanism for the strong scattering reduction of polaritons at high speeds.

DOI: 10.1103/PhysRevLett.108.036405 PACS numbers: 71.36.+c, 42.65.Wi, 67.10.Jn, 78.67.De

The interplay between kinetic energy, localization energy, and particle interactions lies at the heart of the transport properties in condensed matter physics. This problem has been thoroughly treated theoretically since the seminal work of Anderson [1,2], which described the localization of electrons in disordered media caused by the interference between the electronic wave functions.

Electron localization phenomena have been observed in doped semiconductors [3], but this system suffers from a large number of noncontrollable parameters, like the strong Coulomb interactions or the disorder potential. As an alternative, gases of neutral bosons appear as model systems for the understanding of localization. While some works have addressed these questions in the 1980s [4], this activity took an enormous theoretical and experimental expansion since the observation of Bose-Einstein condensates of ultracold atoms [5] and has allowed the study of localization at low dimensions in controlled disorder landscapes. One-dimensional systems are of particular interest as the reduction of the available scattering channels enhances the interference effects. Formally, this gives rise to localization for any disorder strength [6]. Recent experiments have indeed shown the localization of bosonic matter waves in the presence of random disorder using ultracold atoms [7,8].

While disorder tends to localize the wave packets, interparticle repulsive interactions have a delocalization effect. A localized to extended phase transition has been predicted to be triggered by the repulsion between particles in 1D [9,10], and recent experiments have demonstrated this effect in atomic condensates at rest [11]. Particularly interesting is the case of a 1D boson flow in the presence of a weak disordered potential. If the considered sample is small enough and/or the disorder is weak, a quasiextended state is recovered with a high transmission [12], still with an appreciable backscattering amplitude. Repulsive interactions in such a flowing boson gas could lead to the formation of a superfluid state, resulting in the perfect transmission through the 1D channel [9]. However, the superfluid behavior is lost if the flow speed becomes larger than a critical velocity [13] which is close to the speed of sound of the fluid given by $c = \sqrt{\alpha n/m}$, where m is the mass of the particles, n is their density, and α is the interparticle interaction energy. Nevertheless, as we will see in this work, even at high flow speeds in the nonsuperfluid regime, interactions between particles can lead to a strong reduction of the backscattering and, consequently, to an enhancement of the transmission.

In this Letter, we study the effects of interparticle interactions on the propagation of a polariton boson condensate in a 1D semiconductor wire microcavity. We observe, in the supersonic regime, a strong suppression of the signal backscattered by the imperfections of the structure while increasing the condensate density. This suppression, which was previously theoretically predicted in Ref. [9], is caused by the onset of parametric instabilities arising from particle interactions, giving the condensate access to propagating states which bypass the potential barriers caused by disorder. This confirms experimentally the existence of a new mechanism for the suppression of the backscattering, very different from the recently reported superfluidity of polaritons [14,15].

Cavity polaritons are bosonic particles arising from the strong coupling between quantum well excitons and photons confined in a microcavity. Their Bose-Einstein condensation has been reported by several groups [16–19]. Nonlinear phenomena such as optical parametric oscillation (OPO) [20,21], multistable behavior [22], and superfluidity [15,23] have been reported. Recent technological breakthroughs have allowed the production of very high quality microcavities with polariton lifetimes as long as 30 ps—5 times larger than in the preexisting samples. In a 1D geometry, this achievement has allowed the observation of propagating polariton condensates spatially separated from the pumping area and showing phase coherence extending over 100 healing lengths [24]. Here we use this geometry to study the effect of interactions on the backscattering caused by imperfections in the structure.

Our samples are GaAs/AlGaAs-based microcavities etched into 1D wires with a width of 3.5 μm and a length of 0.2 mm [24]. The Rabi splitting amounts to 15 meV and the polariton lifetime to 30 ps. Polariton condensates are formed under nonresonant photoexcitation (laser wavelength around 730 nm) with a single-mode Ti:sapphire laser. The laser spot has a 2 μm diameter and is positioned close to one of the edges of the wire [see Fig. 1(a)]. At low excitation density, carriers are photoinjected at the position of the excitation spot, relaxing down to form polaritons which populate the lower polariton branch. This incoherent gas of polaritons with different momenta can propagate long distances in the wire thanks to the long cavity lifetime. Figure 1(b) shows the momentum space in logarithmic scale of the polariton signal 100 μm away from the injection area. For positive values of the momentum, the forward propagating polaritons can be observed. Because of the large value of the exciton mass, the population of the excitonic reservoir is negligible out of the excitation area.

Above a threshold excitation density P_{th} [Fig. 2(a)], polariton condensation takes place at $k = 0$ in the region optically pumped by the laser. In that region, the presence of a dense excitonic reservoir induces a strong blueshift of the polariton energy even at a few times P_{th}, as discussed in Ref. [24]. The energy difference between the $k = 0$ states *in* and *out* of the pumped area results in an acceleration of the polariton condensate from the pump spot. As a result, in the region of observation, the propagating condensate preserves the same emission energy as in the pumped region but it gains a finite momentum [2.1 μm^{-1} in Fig. 1(c)]. By further increasing the polariton density, the reservoir in the pumped area induces a larger blueshift and the propagating condensate emits at higher energy with a larger k, as shown in Fig. 1(d) and in the black dots of Fig. 1(f).

This progression continues while increasing the excitation density until the propagating condensate gets close to the inflection point of the lower polariton branch. Then, parametric instabilities induced by polariton interactions result in the relaxation of the polariton emission into several coexisting states with different energies [see Fig. 1(e)]. Note that, contrary to the case in nonlinear crystals or in previous OPO experiments in polaritons [20,21], none of the states in the OPO is resonantly fed by a laser source. Here, the OPO is spontaneously triggered once a sufficiently high population is accumulated in the propagating condensate. This is a remarkable example of a self-sustained optical parametric oscillator. As we will see below, these parametric instabilities play an important role in the quenching of the polariton backscattering.

Let us now concentrate on the scattering of the propagating polariton gas with the weak disorder present in the wire microcavity. The origin of the disorder is structural and gives rise to a change of the polariton energy smaller than 0.1–0.4 meV. It is mainly caused by the weak strain accumulated during the growth of the mirrors and also to the nonintentional rugosity that is formed on the walls of the wires during the lithographic and etching processes. Figure 1(b) shows a significant signal backscattered from the incoherent polariton gas at negative momenta. The origin of this signal in the far field can arise only from the scattering of propagating polaritons with disorder as those states cannot be fed by the excitonic reservoir, located 100 μm away from the observation area. When the polariton density is increased above P_{th}, the relative backscattered signal is strongly reduced, as shown in Figs. 1(b)–1(e) and summarized in Fig. 2(b). Simultaneously, the onset of parametric processes is observed in the far field emission (see Fig. 1 and Ref. [25]). The quenching of the backscattering reaches a factor of 20 at the highest investigated excitation density, with respect to the low density regime.

This remarkable reduction is not caused by the increase of the kinetic energy with excitation power, since the

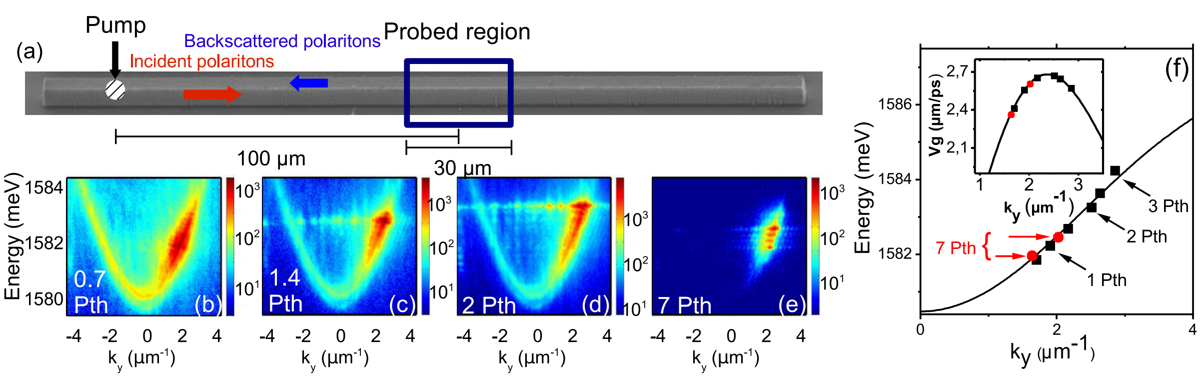

FIG. 1 (color online). (a) Scanning electron microscope image of the wire and scheme of the experiment. (b)–(e) Normalized far field emission of the polariton gas in the probed region. (f) Energy and momentum of the polariton condensates at different excitation powers. The solid line shows the polariton dispersion and the inset the corresponding values of the group velocity.

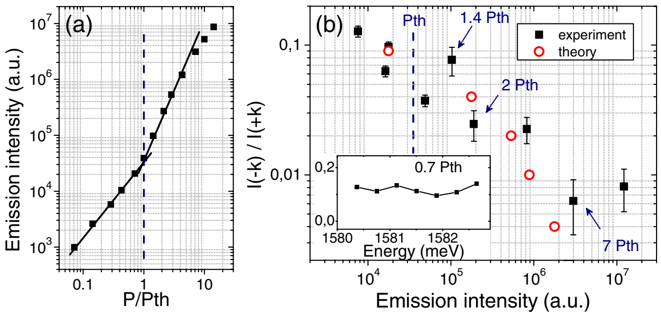

FIG. 2 (color online). (a) Emitted intensity in the probed region as a function of excitation power in units of threshold power P_{th} (7 mW). (b) Black points (red circles): Ratio of the backscattered $[I(-k)]$ to the incident polariton signal $[I(+k)]$ at the energy of the peak signal, experimentally (theoretically) obtained in the conditions of Fig. 1. Inset: Backscattered ratio measured at low excitation power (below the condensation threshold) as a function of the energy of the propagating polaritons.

intrinsic backscattered amplitude, measured at low density, does not depend on the polariton energy [see the inset in Fig. 2(b)]; this intrinsic value is obtained by measuring the backscattered ratio of different polariton states (with different energies) populated along the lower branch below threshold. Note that carrier screening of the disorder is important only in the excitation region, where a dense exciton cloud is present [26]. In the observation area, located far from that spot, any eventual screening of disorder from polariton interactions would arise from a large accumulation of particles in the disorder sites, resulting in bright localized spots in real space and, correspondingly, in a broad emission distribution in reciprocal space, a feature absent in our observations [Figs. 1 and 3].

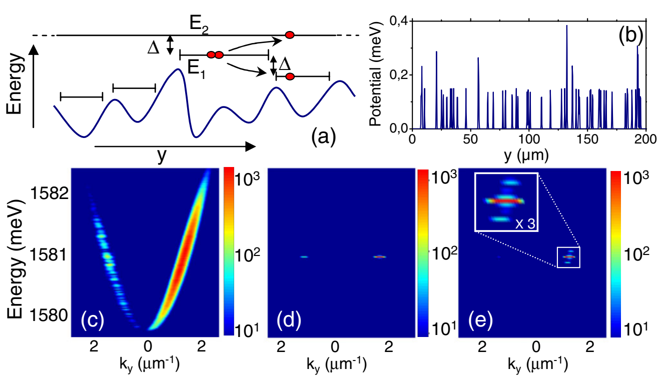

FIG. 3 (color online). (a) Schematic representation of the disorder potential showing the spatial extension of localized states (horizontal bars) and the parametrically induced hopping process. (b) Disorder potential used in the model. (c)–(e) Simulation of the polariton far field emission in the disordered potential with increasing polariton density. The upper-left squared region in (e) shows a magnification of the lower squared area.

The origin of the strong reduction of the scattering can neither be in a transition to a superfluid regime at high density in the context of the Landau criterion, as recently reported in two dimensions under resonant excitation [14,15]. Our condensates travel at large speeds, with momenta on the order of 2–3 $\mu\mathrm{m}^{-1}$ and energy ~ 3 meV above the $k = 0$ state. In order to reach the superfluid regime, interaction-induced blueshifts at least twice the value of the kinetic energy would be required. We are far from that situation, meaning that our fluids are supersonic.

We propose a novel mechanism for the quenching of the backscattering based on the onset of spontaneous parametric processes as the polariton density is increased. In order to understand this mechanism, let us go back to the Anderson model of localization by disorder [1]. One of its simplest descriptions is to consider a regular 1D lattice in the tight-binding model with a hopping constant J, the energy of each site being random [see Fig. 3(a)]. If the energy difference between two sites (E_1 and E_2) is smaller than J ($|E_1 - E_2| < J$), the particle can jump from one site to the other. If the probability of finding an available neighbor for the jump is P (< 1), then the probability to make N jumps will be P^N, which gives an exponentially decaying distribution around the initial site.

Let us consider now a system with local repulsive interaction between particles. This interaction can shift the on-site energy, increasing the probability to match the energies of two different sites. Additionally, the onset of OPO on-site scattering can populate virtually any energy state via the following energy conserving process: $E_1 + E_1 \rightarrow (E_1 + \Delta) + (E_1 - \Delta)$, where 2Δ is the energy difference between signal and idler states. If the condition $|E_1 \pm \Delta - E_2| < J$ is verified, a particle can jump to the neighboring site [Fig. 3(a)].

In our experiment, the kinetic energy of particles is much larger than the localization energy. Therefore, the effect of local blueshift is expected to be weak compared to that of the onset of parametric instabilities, which increases the probability of finding a neighbor with a matching energy and tends to suppress localization. This phenomenon has been theoretically discussed in Ref. [10] in the context of ultracold atoms, while a similar effect has been addressed within a different approach in Ref. [9]. In the latter case, calculations showed that because of the interactions the decay of the transmission of a 1D atomic Bose-Einstein condensate through a disordered waveguide is algebraic instead of exponential, a signature of the suppression of Anderson localization. This regime corresponds to "nonstationary" solutions of the Gross-Pitaevskii equation, evidencing the presence of many-energy states in the system, in line with our interpretation of the onset of parametric processes in experiments.

In order to support our model for the quenching of the backscattering, we have performed numerical simulations of the propagation of polariton condensates with a certain wave vector through a series of pointlike defects. To describe accurately the polariton condensate, we use a set of Schrödinger equations describing the time evolution of the photonic $\psi_{ph}(x,t)$ and excitonic $\psi_{ex}(x,t)$ mean fields coupled via the light matter interaction (Rabi splitting, $\Omega_R = 15$ meV), taking into account the two allowed spin projections of the photon and exciton fields $\sigma = \pm 1$ [27]:

$$i\hbar \frac{\partial \psi^{\sigma}_{ph}}{\partial t} = -\frac{\hbar^2}{2m_{ph}} \Delta \psi^{\sigma}_{ph} + \frac{\Omega_R}{2}\psi^{\sigma}_{ex} - \frac{i\hbar}{2\tau_{ph}}\psi^{\sigma}_{ph} + P^{\sigma} + H_{\rm eff}\psi^{-\sigma}_{ph} + U\psi^{\sigma}_{ph}, \tag{1}$$

$$i\hbar \frac{\partial \psi^{\sigma}_{ex}}{\partial t} = -\frac{\hbar^2}{2m_{ex}} \Delta \psi^{\sigma}_{ex} + \frac{\Omega_R}{2}\psi^{\sigma}_{ph} - \frac{i\hbar}{2\tau_{ex}}\psi^{\sigma}_{ex} + \alpha|\psi^{\sigma}_{ex}|^2\psi^{\sigma}_{ex}. \tag{2}$$

Here $m_{ph} = 3.6 \times 10^{-5} m_0$, $m_{ex} = 0.4 m_0$, and m_0 are the cavity photon, the quantum well exciton, and the free electron masses, respectively. The lifetimes of the particles are $\tau_{ph} = 15$ ps and $\tau_{ex} = 400$ ps. $H_{\rm eff}$ accounts for the TE-TM polarization splitting of the waveguided modes. $\alpha = 6E_b a_B^2/S$ is the polariton-polariton interaction constant ($E_b = 6$ meV is the exciton binding energy, and $a_B = 10$ nm is the exciton Bohr radius), $U(x)$ is the disorder potential along the wire which is a series of randomly spaced delta peaks [shown in Fig. 3(b)], and $P(x,t)$ is the coherent cw pumping, localized on a spot of 2 μm placed on the left of the sample, as sketched in Fig. 1(a), injecting polaritons with well-defined momenta.

Figures 3(c)–3(e) show the far field emission 100 μm away from the excitation spot obtained from the simulations at different excitation densities. To simulate the situation below threshold [Fig. 3(c)], we inject a short pulse, filling all positive momenta states of the lower polariton branch with a distribution similar to the measured one [Fig. 1(b)]. A strong backscattered signal can be observed, analogous to the experimental one shown in Fig. 1(b). Note that the backscattering is entirely due to the presence of the disorder potential. Just above threshold, the condensate presents a single energy [Fig. 1(c)], which we simulate by quasiresonant cw pumping with a well-defined k [Fig. 3(d)]. A significant backscattered signal is still present. At higher densities [Fig. 3(e)], new propagation frequencies associated with the onset of parametric processes become clearly visible, whereas the pumping is still monochromatic. The well-defined separation between the frequencies, which can be observed in both the experimental [Fig. 1(e) and Ref. [25]) and theoretical images [Fig. 3(e)], is due to the polarization splitting $2H_{\rm eff} \approx 0.5$ meV. It determines the preferential resonances for the onset of the OPO processes. This model provides an interpretation to the observation of apparently spectrally equidistant condensates [24] at high pumping rates. The energy spacing between the different frequencies is in our approach not related to a balance between scattering rates and losses [28] but to the value of the TE-TM splitting.

Similar to the results of Ref. [9], the appearance of the new frequencies results in the reduction of the backscattered signal. Figure 2(b) summarizes in open dots the calculated backscattered intensity averaged over 50 disorder realizations, as a function of the excitation density. They show a reduction by more than 1 order of magnitude when the density is increased by a factor of 100, in good agreement with the experimental results. This value is significantly greater than the one found in experiments with condensates at subsonic speeds in planar microcavities (a quenching by a factor of 4) [15]. Though backscattering effects are less important in 2D than in 1D systems due to the very different density of available states in each case, the mechanism we propose might as well play a role in the flow of polaritons in triggered optical parametric configurations in planar microcavities [23,29].

In conclusion, we have observed a strong suppression of the scattering in propagating 1D polariton condensates caused by the onset of parametric instabilities at high excitation densities, which enable the "hopping" of the condensate through the disorder. This new mechanism results in a quasifrictionless flow of supersonic polariton condensates, opening the way to the fabrication of integrated polariton circuits with high transmissivity.

We are greatful to A. V. Kavokin, P. Leboeuf, and N. Pavloff for fruitful discussions. This work was partly supported by the contracts ANR PNANO-07-005 GEMINI, C'Nano Ile de France "Sophiie2," the RTRA "Picorre" and "BoseFlow1D," the ANR-11-BS10-001 "Quandyde," and the FP7 ITNs "Clermont4" (235114) and "Spin-Optronics" (237252).

[1] P. W. Anderson, Phys. Rev. **109**, 1492 (1958).
[2] A. Lagendijk, B. van Tiggelen, and D. S. Wiersma, Phys. Today **62**, No. 8, 24 (2009).
[3] T. F. Rosenbaum, R. F. Milligan, M. A. Paalanen, G. A. Thomas, R. N. Bhatt, and W. Lin, Phys. Rev. B **27**, 7509 (1983).
[4] M. P. A. Fisher, P. B. Weichman, G. Grinstein, and D. S. Fisher, Phys. Rev. B **40**, 546 (1989).
[5] M. H. Anderson, J. R. Ensher, M. R. Matthews, C. E. Wieman, and E. A. Cornell, Science **269**, 198 (1995); K. B. Davis, M. O. Mewes, M. R. Andrews, N. J. van Druten, D. S. Durfee, D. M. Kurn, and W. Ketterle, Phys. Rev. Lett. **75**, 3969 (1995).
[6] N. F. Mott and W. D. Twose, Adv. Phys. **10**, 107 (1961).
[7] J. Billy, V. Josse, Z. Zuo, A. Bernard, B. Hambrecht, P. Lugan, D. Clement, L. Sanchez-Palencia, P. Bouyer, and A. Aspect, Nature (London) **453**, 891 (2008).
[8] G. Roati, C. D'Errico, L. Fallani, M. Fattori, C. Fort, M. Zaccanti, G. Modugno, M. Modugno, and M. Inguscio, Nature (London) **453**, 895 (2008).
[9] T. Paul, P. Leboeuf, N. Pavloff, K. Richter, and P. Schlagheck, Phys. Rev. A **72**, 063621 (2005); T. Paul,

M. Albert, P. Schlagheck, P. Leboeuf, and N. Pavloff, *ibid.* **80**, 033615 (2009).
[10] I. L. Aleiner, B. L. Altshuler, and G. V. Shlyapnikov, Nature Phys. **6**, 900 (2010).
[11] B. Deissler, M. Zaccanti, G. Roati, C. D'Errico, M. Fattori, M. Modugno, G. Modugno, and M. Inguscio, Nature Phys. **6**, 354 (2010).
[12] J. B. Pendry, Adv. Phys. **43**, 461 (1994).
[13] S. Ianeselli, C. Menotti, and A. Smerzi, J. Phys. B **39**, S135 (2006).
[14] I. Carusotto and C. Ciuti, Phys. Rev. Lett. **93**, 166401 (2004).
[15] A. Amo, J. Lefrère, S. Pigeon, C. Adrados, C. Ciuti, I. Carusotto, R. Houdré, E. Giacobino, and A. Bramati, Nature Phys. **5**, 805 (2009).
[16] J. Kasprzak, M. Richard, S. Kundermann, A. Baas, P. Jeambrun, J. M. J. Keeling, F. M. Marchetti, M. H. Szymanska, R. Andre, J. L. Staehli, V. Savona, P. B. Littlewood, B. Deveaud, and L. S. Dang, Nature (London) **443**, 409 (2006).
[17] R. Balili, V. Hartwell, D. Snoke, L. Pfeiffer, and K. West, Science **316**, 1007 (2007).
[18] S. Christopoulos, G. Baldassarri Hoger von Hogersthal, A. J. D. Grundy, P. G. Lagoudakis, A. V. Kavokin, J. J. Baumberg, G. Christmann, R. Butte, E. Feltin, J.-F. Carlin, and N. Grandjean, Phys. Rev. Lett. **98**, 126405 (2007).
[19] E. Wertz, L. Ferrier, D. D. Solnyshkov, P. Senellart, D. Bajoni, A. Miard, A. Lemaitre, G. Malpuech, and J. Bloch, Appl. Phys. Lett. **95**, 051108 (2009).
[20] R. M. Stevenson, V. N. Astratov, M. S. Skolnick, D. M. Whittaker, M. Emam-Ismail, A. I. Tartakovskii, P. G. Savvidis, J. J. Baumberg, and J. S. Roberts, Phys. Rev. Lett. **85**, 3680 (2000).
[21] C. Diederichs, J. Tignon, G. Dasbach, C. Ciuti, A. Lemaître, J. Bloch, P. Roussignol, and C. Delalande, Nature (London) **440**, 904 (2006).
[22] T. K. Paraïso, M. Wouters, Y. Leger, F. Mourier-Genoud, and B. Deveaud-Plédran, Nature Mater. **10**, 80 (2011).
[23] A. Amo, D. Sanvitto, F. P. Laussy, D. Ballarini, E. del Valle, M. D. Martín, A. Lemaitre, J. Bloch, D. N. Krizhanovskii, M. S. Skolnick, C. Tejedor, and L. Viña, Nature (London) **457**, 291 (2009).
[24] E. Wertz, L. Ferrier, D. D. Solnyshkov, R. Johne, D. Sanvitto, A. Lemaitre, I. Sagnes, R. Grousson, A. V. Kavokin, P. Senellart, G. Malpuech, and J. Bloch, Nature Phys. **6**, 860 (2010).
[25] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.108.036405 for spectra showing the onset of parametric processes.
[26] A. Ivanov, L. Smallwood, A. T. Hammack, S. Yang, L. Butov, and A. Gossard, Europhys. Lett. **73**, 920 (2006).
[27] I. A. Shelykh, Y. G. Rubo, G. Malpuech, D. D. Solnyshkov, and A. Kavokin, Phys. Rev. Lett. **97**, 066402 (2006).
[28] M. Wouters, T. C. H. Liew, and V. Savona, Phys. Rev. B **82**, 245315 (2010).
[29] D. Sanvitto, F. M. Marchetti, M. H. Szymanska, G. Tosi, M. Baudisch, F. P. Laussy, D. N. Krizhanovskii, M. S. Skolnick, L. Marrucci, A. Lemaitre, J. Bloch, C. Tejedor, and L. Vina, Nature Phys. **6**, 527 (2010).