Bright solitons in spontaneously formed polariton networks

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Abstract. Bright solitons in a polariton fluid are excitations with a comparatively high intensity that flow without dissipation in spite of a finite lifetime of polaritons. Under resonant driving conditions, bright polariton solitons can be created using a two-beam excitation scheme which includes a plane-wave pump and a focused beam injecting a soliton at a certain location [O. A. Egorov et al., PRL 102, 153904 (2009)]. Here we discuss a new mechanism of soliton formation which takes place under plane-wave driving in a perfectly uniform microcavity wire—i.e., with no writing beams or seeding inhomogeneities. The key phenomenon underlying formation of solitons in this case is the spontaneous breaking of continuous symmetry and onset of periodic spin networks [S. S. Gavrilov, PRL 120, 033901 (2018)].

1. Introduction
Cavity polaritons are composite bosons formed owing to the strong coupling of excitons and cavity photons [1]. They are excited optically and emit light. The lifetimes of polaritons in GaAs-based microcavities lie in the picosecond range, however they form macroscopically coherent (Bose condensate) states under optical driving [2, 3, 4].

In the case of resonant driving, interaction between polaritons results in various collective phenomena, multistability, parametric scattering, etc. Typically, nonequilibrium transitions occur in a threshold manner in response to varying external conditions. Apart from a few number of critical points, polaritons were long thought to follow a resonant driving force adiabatically and, for instance, the state of a homogeneous polariton system driven by a constant plane wave was thought to be also constant. However, coupling between spin components has recently been found to result in nonstationary and, in particular, chaotic polariton states [5].

It is remarkable that turbulence (chaoticity) and long-range ordering in a resonantly driven polariton system are underlaid by exactly the same physics. When all plane-wave states are unstable, the condensate is forbidden to match the symmetry of the external field and its wave vector becomes uncertain [6]. Here the possibilities open up for both chaotic dynamics and the secondary—internal and truly spontaneous—ordering of the system. In particular, a quasi-one-dimensional (1D) microcavity wire arranges itself into a network of spin-up and spin-down domains alternating each other in a strict order [7].

The internally ordered polariton network resembles a crystal rather than a fluid. Similar structures could arise in a periodic potential inducing a lattice of coupled condensates. Today they are often referred to as supersolid states [8, 9], implying a non-dissipative ("superfluid")
propagation of excitations through such a lattice. We have found that polariton networks also exhibit characteristic excitations—solitons—which flow without dissipation and alter their spin states at each node of the network. It is noteworthy, however, that in our case the lattice itself appears out of a perfectly homogeneous Bose gas due to the spontaneous symmetry breakdown.

2. Model and parameters

The model is based on the mean-field approach and very well known. The system evolution is considered in terms of macroscopic wave functions \( \psi_{\pm}(r,t) \) corresponding to the spin-up and spin-down components of the polariton fluid in the microcavity active layer. These functions obey the driven-dissipative Gross-Pitaevskii equations [4],

\[
\begin{align*}
\frac{i\hbar}{\partial t} \psi_+ &= [\hat{E} - i\gamma] \psi_+ + V |\psi_+|^2 \psi_+ + \frac{g}{2} \psi_- + f_+(r,t)e^{-iE_p t/\hbar}, \\
\frac{i\hbar}{\partial t} \psi_- &= [\hat{E} - i\gamma] \psi_- + V |\psi_-|^2 \psi_- + \frac{g}{2} \psi_+ + f_-(r,t)e^{-iE_p t/\hbar}.
\end{align*}
\]

Energy operator \( \hat{E} = \hat{E}(-i\hbar \nabla) \) depends on the dispersion law \( E(k) \); \( \gamma \) is the decay rate (damping coefficient), \( V > 0 \) is the strength of the exchange Coulomb repulsion treated within the dilute-gas approximation, \( g \) is the spin coupling rate, and \( f \) is the driving field amplitude. The pump beam is supposed to be directed at normal incidence, so its in-plane wave vector is zero. Pair interaction between polaritons with opposite spins can be neglected when the system is excited sufficiently below the free exciton level [10]. On the other hand, a strong linear coupling between \( \psi_+ \) and \( \psi_- \) can be provided by a mechanical stress along one of the crystal axes.

We neglect the upper polariton branch that is not excited. The lower-polariton dispersion law reads,

\[
E(k) = \frac{1}{2} [E_C(k) + E_X(k)] - \frac{1}{2} \sqrt{[E_C(k) - E_X(k)]^2 + R^2}.
\]

Here, \( E_{C,X}(k) \) are the energies of the cavity photons \( (C) \) and excitons \( (X) \), which depend on the in-plane wave vector \( k \), and \( R \) is the exciton-photon coupling rate. Cavity photons are characterized by a certain effective mass \( m \) such that \( E_C(k) = E_0 + \hbar^2 k^2 / 2m \). The mass itself is given by the usual formula \( m = \epsilon E_0 / c^2 \), where energy \( E_0 \) and dielectric constant \( \epsilon \) are determined by the microcavity mirrors. The exciton effective mass largely exceeds the photon mass \( m \), thus, the dependence of \( E_X \) on \( k \) can be neglected without reducing generality.

The parameters used for the calculations correspond to a GaAs-based microcavity: \( \epsilon = 12.5, \ E_0 = E_X = 1.5 \text{ eV}, \ R = 10 \text{ meV}, \ g = 0.05 \text{ meV}, \ \gamma = 0.01 \text{ meV} \). The polariton-polariton interaction constant \( V \) was set equal to 1, so that “intensities” \( |\psi_\pm|^2 \) merely yield the blue shifts of the effective resonance energy levels of the two spin components. The pump polarization is linear and perfectly spin-symmetric \( (f_+ = f_-) \). The pump detuning, i.e., the difference between the pump energy level \( E_p \) and the polariton resonance \( E(k=0) \), equals 0.05 meV. The equations are integrated by means of the 4(5) Runge-Kutta method with adaptive step control on a one-dimensional grid representing a microcavity wire with periodic boundary conditions.

In accordance with Ref. [7] the spin and spatial symmetries break down spontaneously in a finite interval of \( f \) when \( g \gtrsim 4\gamma \) and \( g/2 \lesssim D \lesssim 2g \). The chosen parameters meet this criterion.

3. Results and discussion

For the chosen values of the resonance width \( \gamma \), pump detuning \( D \), and spin coupling rate \( g \), several dynamical regimes can be distinguished depending on the pumping amplitude \( f \).

If \( f \) is comparatively small, so that it lies near the beginning of the interval in which all plane-wave solutions of Eqs. (3) are forbidden, then the system arranges itself into a steady and strictly periodic spin network. Increasing amplitude \( f \) results in two different types of excitations: (i)
Figure 1. A fragment of spatiotemporal evolution of a polariton fluid: spontaneously formed spin network and propagation of a bright soliton.

In-situ vibrations of the high-intensity nodes of the network and (ii) bright solitons propagating through the network without losses.

Figure 1 shows the case when the network is relatively stable, but contains a soliton running at the velocity of about 0.3 μm/ps. In the course of propagation, the soliton is seen to alter its circular-polarization degree between +1 and −1. This feature is markedly different from the two-beam scheme in which the polarization of a “manually” created soliton was determined by the excitation conditions [11].

Although a soliton per se is infinitely long-lived, it can collapse by running into network defects or by colliding with other solitons; this is what happens with further increasing amplitude.

Figure 2. Evolution pattern at a larger pump intensity. Left and right panels respectively show the degrees of circular polarization $\rho_{(+,-)}$ and “diagonal” linear polarization $\rho_{(\wedge,\\wedge)}$ (i.e., the third normalized Stokes parameter). The boundary conditions are periodic. The evolution is shown starting from 8 ns in order to exclude transient phenomena.
when the number of irregularities grows significantly. Figure 2 shows the evolution in the case when $f^2$ is 2.5 times greater than in Fig. 1. The observed patterns are no longer perfectly regular, however, they elucidate one important feature of the spin networks. Namely, the spin symmetry can break down in two different ways, so that each pair of neighboring intensity peaks can have either opposite circular polarizations or opposite “diagonal” linear polarizations. Typically, one of the normalized Stokes parameters, $S_1 = \rho(\gamma, \gamma)$ or $S_3 = \rho(\gamma, \gamma)$, is relatively small whereas the other nearly reaches $\pm 1$ at the intensity peaks of the network. The signs of $\rho(\gamma, \gamma)$ and $\rho(\gamma, \gamma)$ are the same at each point. When $f$ is small, the entire network cannot switch between the two ways of symmetry breaking, but such possibility opens up when fluctuations get strong and the network as a whole becomes somewhat uncertain. The solitons seen in the intermediate interval of $f$, which behave as superfluid defects in the lattice of a given type, can also be thought of as “representatives” of the other way of symmetry breaking.

Further increasing $f$ (not shown) makes the spin sites merge into larger domains and—eventually—into a new homogeneous (yet possibly multistable) spin state spreading throughout the system. An essentially different scenario comes into play when $f$, $g/\gamma$ and $D/\gamma$ are increased all together. Here, the sites of the spin network break down into two subsets with synchronized and desynchronized (chaotic) behavior which turn out to be tightly interlaced in space. As a result, order and chaos balance each other and constitute a sort of flowing equilibrium. This class of solutions, known as chimera states, has been specifically considered in Ref. [7].

Finally, let us mention a different route to spin turbulence via multiplying solitons, which takes place under pumping near the magic angle but does not require $g$ be as large as $4\gamma$. In this case, solitons are injected by a focused writing beam, as was done in Ref. [11], however, they show a gradual increase in the number of separate $|\psi_{\pm}|$ peaks, which eventually results in chaotic spin states of polaritons [12].

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References
[1] Weisbuch C, Nishioka M, Ishikawa A and Arakawa Y 1992 Phys. Rev. Lett. 69 3314
[2] Baas A, Karr J P, Romanelli M, Bramati A and Giacobino E 2006 Phys. Rev. Lett. 96(17) 176401
[3] Deng H, Haug H and Yamamoto Y 2010 Rev. Mod. Phys. 82 1489
[4] Kavokin A V, Baumberg J J, Malpuech G and Laussy P 2017 Microcavities 2nd ed (New York: Oxford University Press) ISBN 978-0-19-878299-5
[5] Gavrilov S S 2016 Phys. Rev. B 94(19) 195310
[6] Gavrilov S S 2017 JETP Lett. 105 200–204
[7] Gavrilov S S 2018 Phys. Rev. Lett. 120(3) 033901
[8] Li J R, Lee J, Huang W, Burchesky S, Shteynas B, Top F c, Jamison A O and Ketterle W 2017 Nature 543 91
[9] Léonard J, Morales A, Zupancic P, Esslinger T and Donner T 2017 Nature 543 87
[10] Sekretenko A V, Gavrilov S S and Kulakovskii V D 2013 Phys. Rev. B 88(19) 195302
[11] Sich M, Fras F, Chana J K, Skolnick M S, Krizhanovskii D N, Gorbach A V, Hartley R, Skryabin D V, Gavrilov S S, Cerda-Méndez E A, Biermann K, Hey R and Santos P V 2014 Phys. Rev. Lett. 112(4) 046403
[12] Gavrilov S S 2019 J. Phys. Conf. Ser. 1164 012014