Spinor soliton arrays in cavity-polariton wires

To cite this article: S S Gavrilov 2019 J. Phys.: Conf. Ser. 1164 012014

View the article online for updates and enhancements.
Spinor soliton arrays in cavity-polariton wires

S S Gavrilov
Institute of Solid State Physics, 142432 Chernogolovka, Russia
National Research University Higher School of Economics, 101000 Moscow, Russia
E-mail: gavr_ss@issp.ac.ru

Abstract. Theoretical and numerical study is performed of bright solitons in a resonantly driven cavity-polariton system with linear coupling of two spin components. Together with repulsive Coulomb interaction of polaritons with parallel spins, such coupling involves spontaneous breakdown of spin symmetry. Solitons originate under the joint action of a continuous-wave excitation and spatially local short-term pulse. Due the symmetry breaking, solitons exhibit two intensity peaks with opposite spins which are separated in space. Here we predict spontaneous formation of large soliton arrays with alternating spins. The number of intensity peaks is found to gradually increase in the course of soliton propagation, resulting in turbulent field states.

1. Introduction
Cavity polaritons are composite bosonic excitations originating in semiconductor microcavities in the strong exciton-photon coupling regime [1, 2, 3]. They are very short-lived; for instance, in GaAs-based samples their lifetime is about 10–100 ps. Despite this fact, polaritons can form macroscopically coherent states either due to Bose-Einstein condensation [4, 5], when they are coupled to excitonic reservoir, or under direct resonant and coherent optical driving [6]. In both cases, small perturbations of the condensate usually decay on the scale of the polariton lifetime. A noticeable exception to this rule is given by bright solitons propagating without dissipation through a polariton fluid maintained by a continuous pump wave (cw) whose in-plane wave number exceeds \( k_{\text{infl}} \), the inflection point of the lower-polariton dispersion curve [7, 8, 9]. To create a soliton at a given place, one usually needs to disturb the fluid by a short-term yet sufficiently intensive optical pulse.

Stability of solitons depends on the cw pump amplitude, which has to be set in the interval where the condensate is bistable; then a soliton turns out to be an isolated representative of the high-intensity state on the background of the low-intensity one. In its turn, bistability appears when the pump frequency is slightly blue-detuned with respect to the resonance polariton level; it stems from the repulsive polariton-polariton interaction involving blue shift of eigenfrequencies [10, 11, 12, 13, 14]. If the cw pump is too weak, the state of the condensate is single-valued and stable, so that all fluctuations decay and the solitonic regime is impossible. If, by contrast, the cw pump is too strong, then even a small perturbation triggers a transition of the whole system to the upper steady state. In the latter case, excitation spreads outward the activation point in a neuron-like manner typical of normal-incidence pumping conditions [15, 16]. Variation of the of the short-term pulse parameters has a comparatively lesser effect on the system. In particular, pulses of sufficient duration and/or spatial extent can produce stable...
arrays of solitons running one after another [17]; they are analogous to the soliton trains formerly considered in atomic Bose-Einstein condensates [18].

Here we report a novel solitonic regime: soliton arrays in which the number of intensity peaks with alternating right- and left-circular polarizations gradually grows with time. This kind of solutions is found in a system with a linear coupling of spin components, which implies the splitting of the orthogonally polarized eigenstates due to spatial anisotropy, e.g., thanks to a mechanical stress applied along one of the crystal axes. Systems with such properties were previously investigated under normal-incidence driving; in particular, it was shown that the spin-symmetric states formed when the pump polarization matches the upper eigenstate exhibit spontaneous symmetry breakdown [19, 20, 21, 22]. The circularly polarized solitons discussed in this paper arise because the linearly polarized (spin-symmetric) one-mode state is forbidden. Permanent multiplication of solitons can be thought of as a transition to a new “equilibrium”, which, however, is not a plane-wave state, but rather an internally unstable collective state whose spatial, temporal, and spin symmetries are broken simultaneously.

2. Model

2.1. Equations

Formation of macroscopically coherent polariton states allows one to describe the system within the mean-field approach. The system evolution is considered in terms of macroscopic wave functions $\psi_\pm(\mathbf{r},t)$ of two spin components in the two-dimensional active layer of the microcavity. They obey the driven-dissipative Gross-Pitaevskii equations [3],

$$ i\hbar \frac{\partial \psi_\pm}{\partial t} = \left[ E - i\gamma \right] \psi_\pm + V|\psi_\pm|^2 \psi_\mp + \frac{g}{2} \psi_\pm + f(\mathbf{r},t), $$

(1)

$$ i\hbar \frac{\partial \psi_\mp}{\partial t} = \left[ E - i\gamma \right] \psi_\mp + V|\psi_\mp|^2 \psi_\pm + \frac{g}{2} \psi_\pm + f(\mathbf{r},t). $$

(2)

Here $E = E(-i\hbar \nabla)$ is the energy operator whose explicit form is determined by the dispersion law, $\gamma$ is the decay rate (damping coefficient), $V > 0$ is the polariton-polariton interaction constant, $g$ is the spin coupling rate, and $f$ is the driving field. We assume no pair interaction between polaritons with opposite spins, which is justified in the case when the exciton reservoir is absent [23]. For simplicity, here we consider only the lower polariton branch, which will not affect conclusions. The dispersion law has the form

$$ E(\mathbf{k}) = \frac{1}{2} |E_{C}(\mathbf{k}) + E_{X}(\mathbf{k})| - \frac{1}{2} \sqrt{|E_{C}(\mathbf{k}) + E_{X}(\mathbf{k})|^2 + R^2}, $$

(3)

where $\mathbf{k}$ is an in-plane wave vector, $E_{C,X}(\mathbf{k})$ are the dispersion laws for cavity photons and excitons, correspondingly, and $R$ is the exciton-photon coupling coefficient (the Rabi energy). In a two- or one-dimensional system, photons have a finite effective mass: $E_{C}(\mathbf{k}) = E_0 + \hbar^2 k^2 / 2m$, where $m = \epsilon E_0 / c^2$; energy $E_0$ and dielectric constant $\epsilon$ are determined by the cavity mirrors. The exciton effective mass is approximately 2 orders of magnitude greater than the photon one, so the exciton wave-vector dependence can be neglected [$E_{X}(\mathbf{k}) = \text{const}$].

The external pump is assumed to have the form

$$ f(\mathbf{r},t) = \left( f_1 + f_2 \cdot 2\left( \mathbf{r} - \mathbf{r}_0 \right)^2 / s^2 - (t - t_0)^2 / \tau^2 \right) e^{i(k_p \cdot \mathbf{r} - E_p t / \hbar)}, $$

(4)

where $k_p$ and $E_p / \hbar$ are the in-plane wave vector and frequency, and $f_{1,2}$ are the amplitudes of the background cw pump and short additional perturbation, respectively.
2.2. Parameters

The calculations are performed for a typical GaAs microcavity with the following parameters: \( \epsilon = 12.5, \ E_0 = E_X = 1.5 \ \text{eV}, \ R = 10 \ \text{meV}, \ g = 0.02 \ \text{meV}, \ \gamma = 0.05 \ \text{meV}. \) The polariton-polariton interaction constant was set equal to 1. This choice does not affect the results, but merely fixes the units of \( \psi \) and \( f \); for instance, \( |\psi|^2 \) has the meaning of the blue shift of the resonance energy. The pump wave number \( |\mathbf{k}_p| \approx 1.9 \ \mu\text{m}^{-1} \), which is slightly above the \( E(k) \) inflection point. The pump energy \( E_p \) is 0.4 meV larger than the polariton resonance energy \( E(k_p); f_{1,2} = 3.62 \cdot 10^{-3}. \) The pulse is assumed to arrive at \( t = 0 \) and \( r = 0. \) Its duration and spatial extent are \( \tau = 4 \ \text{ps} \) and \( s = 2 \ \mu\text{m}. \) The equations are solved using the Runge-Kutta method with adaptive step control on a one-dimensional grid representing the microcavity wire.

Figure 1. Top panel: spatiotemporal evolution of the soliton during the first 100 ps; gray color shows the degree of circular polarization, \( \rho = (|\psi_+|^2 - |\psi_-|^2)/(|\psi_+|^2 + |\psi_-|^2). \) Bottom panels: spatial distributions of the \( |\psi_{\pm}| \) components at several reference time moments \( t \); each subplot is centered at \( x - vt \), where \( v \approx 2 \ \mu\text{m}/\text{ps} \) is the soliton velocity.
3. Results and discussion

The calculation results are summarized in the figure. The point \((r = 0, t = 0)\) where the pulse perturbs the steady state is marked by the circle. It is seen that the spin symmetry breaks down in nearly 30 ps after this event, so that the soliton takes the shape of the two-peak pattern with spatially separated spin-up and spin-down components.

In accordance with Eqs. (1)–(2), the spin-up and spin-down components are excited by identical pump sources, so that the model is purely spin-symmetric. It is important that the pump polarization direction matches the upper eigenstate. For instance, using the conventional unitary transformation \(f_{\pm} = (f_x \mp if_y)/\sqrt{2}\) it is easy to see that the pump beam is polarized in the \(x\) direction. On the other hand, the equations are diagonalized at \(|\psi_{\pm}| \to 0\) if one rewrites \(\psi_{\pm}\) via \(\psi_{x,y}\) using the same transformation rule. The eigenenergies \(E_{x,y}\) equal \(E(k_p) \pm g/2\) and, given that \(g > 0\), the \(x\) polarization corresponds to the upper sublevel. As it was shown in Ref.[19], under such conditions the stronger spin component tends to be enhanced further by simultaneously suppressing the minor one, and that is why the spin-symmetric state with high intensity is hardly reachable or is merely unstable with respect to fluctuations. Bearing in mind that a soliton is a representative of the high-intensity state, we see that its formation should be accompanied by the spontaneous spin symmetry breakdown, which is confirmed by the calculations.

Similar effects were experimentally observed in Ref. [24], however, the lifetime of solitons were limited by several tens of picoseconds due to inhomogeneities and finite size of the sample. In a certain parameter range, such two-humped solutions are stable. The solution considered here is obtained at an increased amplitude of the cw pump. It is seen that the soliton is gradually spreading in space upon its propagation through the low-intensity polariton fluid.

On one hand, the obtained solution still demonstrates the solitonic behavior. In particular, it keeps a characteristic velocity of about 2 \(\mu m/ps\) and does not perturb the area it has already gone through; in other words, the velocity of intrinsic spreading is much lower than the velocity of propagation. Second, the solution still exhibits narrow intensity peaks on the micrometer scale because the spin-up and spin-down components compete and alternate each other.

On the other hand, the solution is definitely unstable. In contrast to both a fluctuation decaying with time and a soliton that keeps constant energy, the pattern considered here accumulates energy. This is of no great surprise, because the system is open and driven externally. Moreover, similar behavior was formerly predicted [25] and experimentally observed [26, 27] in the case of normal-incidence driving. Namely, it was shown that the onset of the parametric scattering from the condensate into scattered “signal” and “idler” modes involves—quite counter-intuitively—an increase in the total energy of the system until it reaches its upper steady state. In contrast to previous studies, the “upper” state our system gradually tends to is no longer steady but turbulent in space and time. Analogous solutions (also obtained in the case of normal-incidence driving) were considered in Refs. [28, 29, 30]. The necessary condition for such solutions is that the spin coupling strength \(g\) exceeds the decay rate \(\gamma\) by at least a factor of 4 and the pump intensity is sufficiently strong. Our current results show that turbulence can also occur at lower pump intensities, in which case it manifests itself in growing soliton arrays rather than an immediate transition to chaos in the whole system.

Acknowledgments

The study was supported by the Russian Science Foundation, grant No. 14-12-01372.

References

[1] Weisbuch C, Nishioka M, Ishikawa A and Arakawa Y 1992 Phys. Rev. Lett. 69 3314
[2] Yamamoto Y, Tassone T and Cao H 2000 Semiconductor Cavity Quantum Electrodynamics (Berlin: Springer-Verlag)
[3] Kavokin A V, Baumberg J J, Malpuech G and Laussy P 2017 *Microcavities* 2nd ed (New York: Oxford University Press)

[4] Kasprzak J, Richard M, Kundermann S, Baas A, Jeambrun P, Keeling J M J, Marchetti F M, Szymanska M H, André R, Staehli J L, Savona V, Littlewood P B, Deveaud B and Dang L S 2006 *Nature* **443** 409

[5] Deng H, Haug H and Yamamoto Y 2010 *Rev. Mod. Phys.* **82** 1489

[6] Baas A, Karr J P, Romanelli M, Bramati A and Giacobino E 2006 *Phys. Rev. Lett.* **96**(17) 176401

[7] Egorov O A, Skryabin D V, Yulin A V and Lederer F 2009 *Phys. Rev. Lett.* **102**(15) 153904

[8] Egorov O A, Skryabin D V and Lederer F 2011 *Phys. Rev. B* **84**(16) 165305

[9] Sich M, Krizhanovskii D N, Skolnick M S, Gorbach A V, Hartley R, Skryabin D V, Cerda-Mendez E A, Biermann K, Hey R and Santos P V 2012 *Nat. Photon.* **6** 50

[10] Baas A, Karr J P, Eleuch H and Giacobino E 2004 *Phys. Rev. A* **69** 023809

[11] Gippius N A, Tikhodeev S G, Kulakovskii D V, Krizhanovskii D N and Tartakovskii A I 2004 *Europhys. Lett.* **67** 997

[12] Gavrilov S S, Gippius N A, Kulakovskii V D and Tikhodeev S G 2007 *JETP* **104** 715

[13] Gippius N A, Shelykh I A, Solnyshkov D D, Gavrilov S S, Rubo Y G, Kavokin A V, Tikhodeev S G and Malpuech G 2007 *Phys. Rev. Lett.* **98** 236401

[14] Gavrilov S S, Gippius N A, Tikhodeev S G and Kulakovskii V D 2010 *JETP* **110** 825

[15] Liew T C H, Kavokin A V and Shelykh I A 2008 *Phys. Rev. Lett.* **101** 016402

[16] Amo A, Liew T C H, Adrados C, Houdré R, Giacobino E, Kavokin A V and Bramati A 2010 *Nat. Photon.* **4** 361

[17] Chana J K, Sich M, Fras F, Gorbach A V, Skryabin D V, Cancellieri E, Cerda-Méndez E A, Biermann K, Hey R, Santos P V, Skolnick M S and Krizhanovskii D N 2015 *Phys. Rev. Lett.* **115**(25) 256401

[18] Strecker K E, Partridge G B, Truscott A G and Hulet R G 2002 *Nature* **417** 150–153

[19] Gavrilov S S, Sekretenko A V, Novikov S I, Schneider C, Höfling S, Kamp M, Forchel A and Kulakovskii V D 2013 *Appl. Phys. Lett.* **102** 011104

[20] Sekretenko A V, Gavrilov S S, Novikov S I, Kulakovskii V D, Höfling S, Schneider C, Kamp M and Forchel A 2013 *Phys. Rev. B* **88**(20) 205302

[21] Gavrilov S S, Sekretenko A V, Gippius N A, Schneider C, Höfling S, Kamp M, Forchel A and Kulakovskii V D 2013 *Phys. Rev. B* **87**(20) 201303

[22] Gavrilov S S, Brichkin A S, Novikov S I, Höfling S, Schneider C, Kamp M, Forchel A and Kulakovskii V D 2014 *Phys. Rev. B* **90**(23) 235309

[23] Sekretenko A V, Gavrilov S S and Kulakovskii V D 2013 *Phys. Rev. B* **88**(20) 195302

[24] 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

[25] Gavrilov S S 2014 *Phys. Rev. B* **90**(20) 205303

[26] Gavrilov S S, Brichkin A S, Grishina Y V, Schneider C, Höfling S and Kulakovskii V D 2015 *Phys. Rev. B* **92**(20) 205312

[27] Whittaker C E, Dzurnak B, Egorov O A, Buonaiuto G, Walker P M, Cancellieri E, Whittaker D M, Clarke E, Gavrilov S S, Skolnick M S and Krizhanovskii D N 2017 *Phys. Rev. X* **7**(3) 031033

[28] Gavrilov S S 2016 *Phys. Rev. B* **94**(19) 195310

[29] Gavrilov S S 2017 *JETP Lett.* **105** 200–204

[30] Gavrilov S S 2018 *Phys. Rev. Lett.* **120**(3) 033901