Magnetic control over the zitterbewegung of exciton–polaritons

E S Sedov1,2,3, I E Sedova3, S M Arakelian3 and A V Kavokin1,2,4

1 School of Science, Westlake University, 18 Shilongshan Road, Hangzhou 310024, Zhejiang Province, People’s Republic of China
2 Institute of Natural Sciences, Westlake Institute for Advanced Study, 18 Shilongshan Road, Hangzhou 310024, Zhejiang Province, People’s Republic of China
3 Vladimir State University named after A. G. and N. G. Stoletovs, Gorky str. 87, 600000, Vladimir, Russia
4 Spin Optics Laboratory, St-Petersburg State University, 1 Universitetskaya, St-Petersburg 198504, Russia

E-mail: evgeny.sedov@mail.ru

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Abstract
The effect of the zitterbewegung consisting in trembling of trajectory of propagating particles may, in principle, be found in a variety of physical systems characterized by split kinetic energy dispersion branches. However, in a majority of material systems the effect is too weak to be observable. Specially designed semiconductor heterostructures representing optical microcavities with embedded quantum wells allow observing the zitterbewegung of exciton–polaritons that are optical cavity modes strongly hybridized with excitons in quantum wells. Here we show that external magnetic fields applied in the plane of the microcavity amplify this effect and allow for tuning the amplitude and the period of oscillations of polariton trajectories, thus being a convenient tool of control. These results pave the way towards realization of ballistic polariton transistors based on the spin–orbit effect, conceptually similar to Datta-and-Das transistors.

1. Introduction
The effect of the zitterbewegung manifests itself in the appearance of oscillations of trajectories of propagating particles due to the correlations between some of its degrees of freedom. First predicted for free Dirac electrons [1], it was later was generalized to a wide variety of systems characterized by split kinetic energy dispersion branches. Among them are trapped ions [2], Bose–Einstein condensates of ultracold atoms [3], two-dimensional photonic crystals [4], binary waveguide lattices [5], mono- and bilayer graphene, and carbon nanotubes [6]. A separate group of systems with gapped kinetic energy spectra are spin-split systems, where the degeneracy in energy is lifted for states with different internal degrees of freedom (spin and/or polarization). The spin–orbit coupling (SOC) which is the mutual influence of the internal degree of freedom of a particle and the external degrees of freedom associated with its macroscopic behavior [7, 8], is at the origin of the characteristic energy splittings of kinetic energy branches in such systems. The manifestation of the zitterbewegung oscillations induced by the spin–orbit interaction (SOI) in a condensate of ultracold boson atoms was reported in [9]. In solid state systems, the zitterbewegung of electrons induced by the SOC due to the Rashba and Dresselhaus effects in semiconductor structures was demonstrated [10, 11].

In the recent paper [12], some of us predicted theoretically the zitterbewegung of macroscopic wave packets of exciton–polaritons in a two-dimensional optical microcavity with embedded quantum wells (QWs). Exciton–polaritons are mixed states formed from the microcavity photon modes strongly coupled to the QW excitons [13]. The predicted in [12] zitterbewegung comes from the optical constituent of the polariton state. The splitting in the transverse electric and transverse magnetic cavity modes (TE–TM splitting) is at the origin of this effect. The magnitude and observability of the predicted zitterbewegung are limited by the fact that the period of oscillations increases quadratically with the polariton wave number $k$, and because the amplitude of the oscillations is inversely proportional to $k$. Moreover, the contribution of
the photonic component to the exciton–polariton state weakens with the increasing $k$. Consequently, the effect is only observable in a narrow range of $k$ close to the bottom of the low-branch polariton dispersion, where it closely approaches the parabolic cavity photon dispersion. This also results in a fast broadening of the polariton wave packet which obscures the effect of the zitterbewegung.

In the present work, we are taking advantage of the composite nature of exciton–polaritons that allows one to significantly expand the conditions favorable for the manifestation of the zitterbewegung. Taking into account the energy splitting of the exciton state induced by the external magnetic field we were able to efficiently compensate reduction of the zitterbewegung caused by the weakening of the photon contribution with the increasing wave vector. In [14–18], it was shown that the splitting of linearly polarized exciton states occurs under the effect of the external magnetic field applied in the Voigt geometry (in the cavity plane) as a result of the magneto-induced mixing of optically active (bright) exciton states with optically inactive (dark) states. The magnitude of the splitting depends quadratically on the magnitude of the applied magnetic field. The external magnetic field is known as an effective tool for controlling polarization properties of exciton–polaritons, however the Voigt geometry remains undeservedly deprived of attention, being overshadowed by the Faraday geometry [19–23]. The recent paper [24] is devoted to the experimental study of the joint effect of the TE–TM splitting of cavity modes and the splitting of excitons induced by the in-plane magnetic field on the polariton polarization dynamics in real space. The variation of the period of polarization beats due to the interplay of photon and exciton mechanisms of the energy splitting and the suppression of oscillations when the splitting effects compensate each other were demonstrated.

In the discussed experiments, as well as in the vast majority of the previous studies of the spin-split excitons, attention was paid to the effect of SOI on spin (polarization) transport of excitons or exciton–polaritons, see e.g. [22–31]. Here we underline that SOI implies the mutual influence of the spin and orbital degrees of freedom of polaritons. In this manuscript, we focus in particular on the spin-to-orbital influence and study the effect of polarization on the trajectories of ballistically propagating polaritons. We show that the contribution of the magneto-induced splitting of the exciton constituent allows polaritons to exhibit the zitterbewegung oscillations at large wave vectors, even above the polariton dispersion inflection point. The suppressed broadening of the polariton wave packets in the region of the flatter exciton-like dispersion makes the oscillations apparent over longer distances. The visibility and sustainability of oscillations are improved due to the optical pumping of the trembling polariton state. The real-space splitting of the polariton wave packet is demonstrated as a result of the interplay between the TE–TM splitting of the cavity modes and the magneto-induced anisotropic splitting of excitons.

2. An analytical background of the polariton zitterbewegung effect

The geometry of the possible experiment schematically shown in figure 1(a) conceptually coincides with that suggested in [12]. Exciton–polaritons are created by a resonant cw optical pump in an optical microcavity with embedded QWs operating in a strong-coupling regime. To increase the propagation distance of the polariton wave packet, we add the spatially-homogeneous non-resonant optical pump. The main difference between two schemes is the applied in-plane magnetic field.

The Hamiltonian describing the behavior of spin-split exciton–polaritons of the lower branch in the microcavity plane can be represented as

$$\hat{H}_0 = \hat{T} + \Omega \cdot \hat{S}. \quad (1)$$

It is important to emphasize that in our case, in contrast to [25], the polariton kinetic energy described by the operator $\hat{T}$ is anisotropic and accounts for the non-parabolicity of the lowest polariton branch:

$$\hat{T} = \frac{1}{2} \left[ \hat{s} - \hat{E}_R \right] \sigma_0, \quad (2)$$

where $\hat{s} = \Delta + \hbar^2 k^2 / 2 m_C$ describes the $k$-dependent exciton–photon detuning. We neglect a dependence on $k$ of the energy of excitons and choose it as a reference. $\Delta$ is a constant detuning component, $m_C$ is the effective mass of cavity photons, $k = (k_x, k_y)$ is the polariton quasimomentum operator, $\hat{E}_R = (\hat{\delta}^2 + V_R^2)^{1/2}$ is the $k$-dependent splitting of the polariton branches, $V_R/2$ is the vacuum Rabi energy. $\sigma_0$ is the $2 \times 2$ identity matrix.

We describe the polariton quantum state with the spinor wave function $|\Psi\rangle = [\Psi_+ (r), \Psi_- (r)]^T$, where $\Psi_{\pm (r)} (r)$ characterize right-circularly (left-circularly) polarized polariton states. We introduce the three-dimensional pseudospin operator $\hat{S} = \frac{1}{2} \hat{\sigma}$, where $\hat{\sigma} = (\sigma_x, \sigma_y, \sigma_z)$ is the vector of the Pauli operators. The second term in the Hamiltonian (1) is responsible for SOI of polaritons. The contributions from both the photon and exciton fractions give rise to the effective magnetic field $\Omega = \Omega_C + \Omega_X$, which causes precession of the exciton pseudospin. The effective field $\Omega_C$ comes from the cavity photon fraction of the
polaron state and originates from the TE–TM splitting of the cavity modes: 
\[ \Omega_C = \begin{bmatrix} \beta C_2^2(k_x^2 - k_y^2), 2\beta C_2^2k_xk_y, 0 \end{bmatrix}, \]
where \( \beta \) is the splitting constant, \( C_2^2 = (1 - \delta/E_0)/2 \) is the Hopfield coefficient which determines the photon fraction in the polaron state.

Another component \( \Omega_X \) of the effective magnetic field \( \Omega \) determined as \( \Omega_X = [\eta X_2^2(B_x^2 - B_y^2), 2\eta X_2^2B_xB_y, 0] \) describes the contribution of the external in-plane magnetic field \( B = (B_x, B_y, 0) \) affecting the exciton fraction of polaritons. The splitting constant \( \eta \) results from magnetic-field-induced mixing of bright and dark exciton states and is a characteristic of the considered structure [16–18, 24]. The weighting Hopfield coefficient for the exciton fraction is \( X_2^2 = 1 \) and \( X_2^2 + X_4^2 = 1 \).

Using the generalized commutation relations, we arrive at the equations for the trajectory operators
\[ \partial_t \hat{x} = \hbar^{-1} \partial_{k_x} \hat{H}_0, \quad \partial_t \hat{y} = \hbar^{-1} \partial_{k_y} \hat{H}_0. \]

The precession equation for the pseudospin operator preserves its form \( \hbar d\hat{S} = \hat{\Omega} \times \hat{S} \) given in [12, 25] regardless the form of the effective magnetic field operator \( \hat{\Omega} \).

We take the polaron wave function as \( |\Psi\rangle = \Psi(r) |\psi\rangle \), where the function \( \Psi(r) = \hat{S}^{-1}[\Psi(k - k_0)] \) is responsible for the spatial envelope, and the spinor \( |\psi\rangle = [\psi_+, \psi_-]^T \) describes the polarization state. \( \hat{S}^{-1} \) is the inverse Fourier transform operator. To reveal the \textit{zitterbewegung} of the polaron wave packet, we now take the same particular conditions as in [12]. We consider the circularly polarized wave packet, \( |\psi\rangle = (1, 0)^T \), propagating ballistically in \( y \) direction, \( k_0 = (0, k_0) \). We also take the external magnetic field applied in \( x \) direction, \( B \uparrow \uparrow x \), i.e. we assume \( B_x = B, B_y = 0 \). In the quasi-classical limit applicable, where the spectrum of the wave packet \( \Psi(k - k_0) \) centered at \( k_0 \) is narrow, we obtain the trajectory
\[ X(t), Y(t) = [(\hat{x}), (\hat{y})] \] of the polaron wave packet in the form:
\[ X(t) = -\frac{\Omega_0}{\Omega_{0\delta}} \left[ 1 - \cos \left( \hbar^{-1} \Omega_{0\delta} t \right) \right], \quad Y(t) = \frac{\hbar k_0}{m^*} t, \] while the polarization vector evolves as \( S(t) = (\hat{S}) = \left[ 0, -\sin(\hbar^{-1} \Omega_{0\delta} t), \cos(\hbar^{-1} \Omega_{0\delta} t) \right]^T \). \( \Omega_{0\delta}, \Omega_0 \) characterize the components of the effective magnetic field \( \Omega_0 \) affecting the polaron propagation. The latter can be obtained from \( \hat{\Omega} \) by substituting \( k \to k_0 \). Thence
\[ \Omega_{0\delta} = \Omega_x|k = k_0| = -\beta C_2^2 k_0^2 + \eta X_2^2 B_x^2, \]
\[ \Omega_0 = (\partial_{k_x} \Omega_y)|k = k_0| = 2\beta C_2^2 k_0 k_y. \]
The discussion above is valid for a ballistically propagating polariton in the quasi-classical limit. The numerical model for simulating the zitterbewegung is described in the next section. As one can see, there is a region on the phase plane $k_0 – B$, when the amplitude $A$ of the zitterbewegung is considerable and, in theory, can be infinitely large. It happens when the TE–TM splitting and the anisotropic splitting are of the same order and tend to compensate each other, so the component of the effective field $\hat{B}$ happens when the TE–TM splitting and the anisotropic splitting are of the same order and tend to compensate each other, so the component of the effective field $\Omega_0$ (5a) tends to zero. The dash-dotted curves in figure 2 denote the parametric dependence of the parameters $B_1(k_0)$ for which $\Omega_0 = 0$. At the same time, with the increase of $A$, the period of oscillations $L \propto \Omega_0$ increases proportionally and tends to the singularity in the same conditions, see figure 2(b). Conversely, when the period $L$ is small enough for observing the zitterbewegung, e.g. at $k_0$ close to zero, its amplitude is negligibly small, which makes the oscillations indistinguishable. Nevertheless, there exist regions in the phase plane with the parameters, for which the amplitude of the zitterbewegung remains considerably large (of several micrometers) at reasonably small periods (several hundreds of micrometers).

3. The numerical model for simulating the zitterbewegung

The discussion above is valid for a ballistically propagating polariton in the $k_0$ state. However, the finiteness of the spectrum of a polariton wave packet as well as non-conservative processes accompanying the polariton evolution significantly affect the phase diagrams on the phase plane $k_0 – B$. To describe the zitterbewegung in experimentally realizable conditions, we numerically solve the generalized Pauli equation in the basis of right- and left-circular polarizations, $|\Psi\rangle = \begin{bmatrix} |\Psi_+\rangle(t, \mathbf{r}) & |\Psi_-\rangle(t, \mathbf{r}) \end{bmatrix}^T$:

$$i\hbar \partial_t |\Psi\rangle = \begin{bmatrix} H_0 + i\hbar \frac{1}{2}(\hat{B} - \gamma^c) \end{bmatrix} |\Psi\rangle + i |F\rangle,$$

where we take into account non-conservative gain and loss processes. The operator $\gamma^c = \gamma^c_{\omega_c} + \gamma^x_{\omega_x}$ is responsible for the polariton decay rate which in general case is $k$-dependent, $\gamma^c_{\omega_c}$ and $\gamma^x_{\omega_x}$ are the decay rates of cavity photons and QW excitons. Polaritons are created by the resonant probe $|F\rangle = F(\mathbf{r})|f\rangle$, where the function $F(\mathbf{r})$ is taken in the Gaussian form $F(\mathbf{r}) \propto \exp \left[-\frac{1}{2}(2\pi\sigma_p)^2\right] \exp\{i(k_p r - \omega_p t)\}$, $k_p$ and $\omega_p$ are the wave vector and the frequency of the probe, respectively, $\sigma_p$ is the width of the probe spot. In further

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**Figure 2.** Characteristics of the zitterbewegung of exciton–polaritons in the quasi-classical limit. (a) The amplitude $A$ and (b) the period $L$ of the zitterbewegung as functions of the magnitude of the external magnetic field $B$ and the wave number $k_0$ of polariton propagation estimated in the quasi-classical limit from (4). The cavity TE–TM splitting constant is taken as $\beta = 140 \mu eV \mu m^2$. The other parameters are the same as in figure 1. The color schemes for the panels (a) and (b) are chosen such that the preferable values of the parameters (larger $A$ and smaller $L$) belong to the red spectral range. The dash-dotted curves (black in (a) and white in (b)) indicate the singularity in the parameters achieved at $\Omega_0 = \Omega_X$ for the $k_0$ polariton state. The gray equivalence contours on the color maps are the guide for the eye.
simulations, we take \( k_p \) and \( \omega_p \) not independently, but matching the polariton dispersion. The vector \(|f\rangle = (f_+, f_-)|^3\) characterizes the polarization of the probe.

To support formation of the polariton condensate, we introduce the non-resonant optical pumping of polaritons via the reservoir of excitons [12, 32]. The operator \( \hat{R} = \hat{R}[n_{k\pm}(t, r)(\hat{\sigma}_0 + \hat{\sigma}_z) + n_{k\mp}(t, r)(\hat{\sigma}_0 - \hat{\sigma}_z)]/2 \) describes the inflow from the reservoir to the polariton state due to the stimulated scattering with the rate \( n_{k\pm}(t, r) \) are the densities of the reservoir excitons of the corresponding polarizations. The evolution of the reservoir described by the spinor \(|n_R\rangle = [n_{k+}(t, r), n_{k-}(t, r)]\) obeys the rate equation

\[
\frac{\partial |n_R\rangle}{\partial t} = (\hat{P} - (\gamma_X + \hat{W})|n_R\rangle,
\]

where the operator \( \hat{W} \) describing the outflow from the reservoir can be obtained from \( \hat{R} \) by replacing \( n_{k\pm}(t, r) |\Psi_{k\pm}(t, r)|^2 \) describes the non-resonant pumping of the exciton reservoir. The pump is considered spatially homogeneous and working in the so-called ‘dark regime’, when the pump power does not exceed the threshold power \( P_{th} \). The polariton wave packet cross-section preserves its shape and width allowing to observe the remarkable peculiarity of all the plots is a much weaker spreading of the polariton wave packet in comparison with one discussed in [12] due to approaching the flatter exciton-like region of the polariton dispersion. The polariton wave packet cross-section preserves its shape and width allowing to observe the

4. The zitterbewegung of polariton wave packets

To trace the effect of the external in-plane magnetic field of the zitterbewegung of polariton wave packets in experimentally realizable conditions, we perform a set of 1600 numerical experiments with different values of the parameters \( k_p \) and \( B \). Here the wave number of the probe \( k_p \) is used instead of \( k_0 \). The width of the probe spot is taken as \( w_p = 10 \mu m \). The TE–TM splitting constant is taken as \( \beta = 140 \mu eV \mu m^2 \). Figure 3 summarizes the results of the simulations and shows the variation of the amplitude \( A \) (figure 3(a)) and the period \( L \) (figure 3(b)) of the zitterbewegung in the phase plane \( k_p - B \). Since the amplitude of oscillations of the trajectory of the polariton wave packet of a finite width decays with the increase of the distance from the injection spot, for the amplitude \( A \) of the zitterbewegung, we take the maximum deviation of the trajectory in \( x \) direction on the first period of oscillations. For the ease of comparison, the ranges of values of \( k_p \) and \( B \) are taken the same as those in figure 2. Due to the limitations of our numerical experiment caused by the finite size of the computational grid, we were unable to estimate the parameters \( A \) and \( L \) in the case where the period \( L \) was large compared to the size of the calculation area. We shaded the corresponding regions in white (in figure 3(a)) an dark blue (in figure 3(b)). A very good qualitative agreement in figures 2 and 3 shows that the main conclusions made for the quasi-classical limit apply to real polariton wave packets. For the period \( L \) of the zitterbewegung, even the quantitative agreement is remarkable. Far enough from the singularity curve, the discrepancy in the numerical estimations and the analytical predictions hardly exceeded 10%. As for the amplitude \( A \) of the zitterbewegung, according to figure 3, it is considerably smaller than that in the quasi-classical limit. The discrepancy reaches tens of percent and it is the larger, the closer the parameters \( (k_p, B) \) are to the critical values \( (k_c, B_c) \).

As one can see from figure 3, under the in-plane magnetic field, the zitterbewegung becomes observable at large wave numbers even above the inflection point of the polariton dispersion. For the used parameters, for polaritons propagating in \( y \) direction, the dispersion possesses the inflection point at \( k_{inf} \approx 1.61 \mu m^{-1} \) and \( \hbar \omega_{inf} \approx -1.7 \text{meV} \). To illustrate the effect of the external magnetic field on the zitterbewegung, in figures 4(a)-(f), we show the spatial distribution of the density of the polariton condensates \( I = \Psi^\dagger \Psi \) in the steady state at different values of the magnetic field magnitude \( B \). We take the probe such that it creates polaritons above the inflection point, in particular, \( k_p = 2 \mu m^{-1} \) and \( \hbar \omega_p \approx -1.4 \text{meV} \). The most remarkable peculiarity of all the plots is a much weaker spreading of the polariton wave packet in comparison with one discussed in [12] due to approaching the flatter exciton-like region of the polariton dispersion. The polariton wave packet cross-section preserves its shape and width allowing to observe the
Figure 3. Characteristics of the zitterbewegung of exciton–polaritons in experimentally realizable conditions. (a) The amplitude $A$ and (b) the period $L$ of the zitterbewegung as functions of the magnitude of the external magnetic field $B$ and the wave number $k_p$ of the probe estimated from the numerical simulations of (6). The width of the spot of the probe is taken $w_p = 10 \, \mu m$. Each square on the color maps corresponds to a single numerical experiment. The squares are colored in white in (a) and dark blue in (b) for the numerical experiments, when we were unable to estimate values of the parameter due to calculation limitations. Values of the parameters used for simulations are given in the text.

Figure 4. The effect of the external in-plane magnetic field on the zitterbewegung of exciton polaritons. (a)–(f) The effect of the zitterbewegung on the spatial distribution of the density the polariton condensate at different values of the magnitude of the external magnetic field (given in the panels). The polariton condensates were created by the probe beam of width of $w_p = 10 \, \mu m$ with the wave number of $k_p = 2 \, \mu m^{-1}$. (g) The trajectories $X(Y)$ of the centers of mass of the polariton wave packets shown in the panels (a)–(f). (h) The amplitude $A$ and (i) the period $L$ of the zitterbewegung of the polariton wave packet at $k_p = 2 \, \mu m^{-1}$ as functions of $B$. The color markers on the panels (a)–(f) and (h)–(i) indicate the curves of the corresponding color on the panel (g).

zitterbewegung at much larger distances. Figure 4(g) shows the corresponding weighted average trajectories of polaritons $X(Y) = \langle \Psi | x | \Psi \rangle_x / \langle \Psi | \Psi \rangle_x$, where $\langle \cdots \rangle_x$ denotes averaging over $x$ in the steady state. For the convenience of analysis, in figures 4(h) and (i) we show the dependencies of the amplitude $A$ and the period $L$ of the zitterbewegung on $B$ at a fixed value of $k_p = 2 \, \mu m^{-1}$. In the considered conditions, the TE–TM splitting matches the magneto-induced splitting at $B_c \approx 13.28 \, T$. Figures 4(g)–(i) confirm the simultaneous increase in the amplitude $A$ and the period $L$ of the zitterbewegung with $B$ approaching $B_c$. Herewith, the apparently more steep curve in figure 4(i) as compared to figure 4(h) shows that the increase in $L$ is faster than one in $A$. 
We would like to point out that the phase of the zitterbewegung (which is responsible for the direction of zitterbewegung of the polariton wave packet of the finite size) is opposite for $B < B_c$ and $B > B_c$. This is connected to the different orientations of the effective magnetic field $\mathbf{Ω}(\mathbf{k}_p) = (\Omega_{\varphi}, 0, 0)$ at different $B$ values which is determined by the sign of the $x$ component $\Omega_{\varphi} \equiv \Omega_x|k_x-k_p|$. If the external magnetic field magnitude $B$ is smaller than the critical value $B_c$, the contribution from the TE–TM splitting exceeds one from the magneto-induced splitting, $|\Omega_{\text{C},\chi}| > |\Omega_{\text{X},\chi}|$ and $\Omega_{\varphi} < 0$, which results in the effective magnetic field oriented opposite to the direction of polariton wave packet at $k_p \equiv 2 \mu m^{-1}$ and $B = 12 \, T$ as a function of $w_p$. The color markers in (b) indicate the curves of the corresponding color in (a).

When the magnetic field magnitude $B$ is close to $B_c$, the spatial polariton density distribution deviates from the solid shape and acquires a fork-like shape. To clarify this, in figure 5 we show the spatial distribution of the pseudovector components $\mathbf{S}_j = \Psi^\dagger(\mathbf{r})\hat{S}_j\Psi(\mathbf{r})/I(\mathbf{r})$ ($j = x, y, z$) characterizing polarization of the exciton–polariton condensates of the solid shape at $B = 10 \, T$ (figure 5(a)) and of the fork-like shape at $B = 13 \, T$ (figure 5(c)). At $B = 10 \, T$, the pronounced regular patterns are characteristic to the $S_x$ component supplemented by wavy patterns in the $S_z$ component. The period of oscillations of the polarization patterns coincides with one of the zitterbewegung. At $B = 13 \, T$, the situation dramatically changes. The splitting in linear polarizations is small which causes mixing of the linearily polarized modes. The polarization plane of the polariton states rotates as a result of this mixing, herewith the direction of rotation is $k_x$-dependent. In figure 5(c) on the color map for the diagonal component $S_x$, instead of oscillations in $y$ direction we observe the separation in $x$ direction of the $S_x > 0$ and $S_x < 0$ polarizations reflected in the fork-like shape of the polariton density distribution. With $B$ approaching the critical value $B_c$, the separation of opposite diagonal polarizations becomes more pronounced, while the contribution of $S_{y}$ and $S_{z}$ diminishes. The color maps of polarization components for all polariton condensates shown in figure 4 are given in the supplementary information.

As we mentioned earlier, the amplitude of the zitterbewegung of the polariton wave packet of the finite width is smaller than one in a quasi-classical limit. Figure 6 illustrates the effect of width of the probe beam $w_p$ on the zitterbewegung of polaritons. Center-of-mass trajectories $X(Y)$ of polaritons are plotted for different values of the width of the probe beam: $w_p = 5 \, \mu m$ (green curve), $w_p = 10 \, \mu m$ (blue curve) and $w_p = 25 \, \mu m$ (crimson curve). (b) The amplitude $A$ of the zitterbewegung of the polariton wave packet at $k_p \equiv 2 \mu m^{-1}$ and $B = 12 \, T$ as a function of $w_p$. The color markers in (b) indicate the curves of the corresponding color in (a).
demonstrated in [39]. The heavy-light hole mixing can also been provoked by an external impact. In [40],

stage of growth, the built-in electric field of roughness of interfaces [36–38]. The magnitude of the splitting
mixed states. The symmetry breaking can be inherent to the structure due to, e.g. an acquired strain at the

observation of this important effect for polaritons with large wave vectors even above the inflection point.

photon polarization splitting that appears in the presence of the in-plane magnetic field makes possible the

below the inflection point of the low-polariton dispersion branch [12]. The interplay between exciton and

leads to the appearance of an internal electric field. This field mixes heavy holes, which are in the basis of

anisotropy [34, 35], the reduced symmetry of the QW interfaces in comparison with the bulk crystal that

optically excited polaritons, and light holes, herewiththe splitting in energy is characteristic to the new

scattering of excitons from the reservoir to the polariton mode [32]. The reservoir is conveniently created by

these significant propagation distances are characteristic of photon-like polaritons. In our manuscript, only the magnetic-field-induced splitting of excitons was considered. Among other possible mechanisms of the splitting are the exchange induced anisotropy [34, 35], the reduced symmetry of the QW interfaces in comparison with the bulk crystal that leads to the appearance of an internal electric field. This field mixes heavy holes, which are in the basis of optically excited polaritons, and light holes, herewith the splitting in energy is characteristic to the new mixed states. The symmetry breaking can be inherent to the structure due to, e.g. an acquired strain at the stage of growth, the built-in electric field of roughness of interfaces [36–38]. The magnitude of the splitting in this case as a rule amounts of tens of microelectron volt. The strain-induced splitting of 100 μeV was demonstrated in [39]. The heavy–light hole mixing can also been provoked by an external impact. In [40], the authors apply stress to the sample and achieve impressive values of the splitting of exciton states of up to 1 meV in specific conditions. Nevertheless, the magnetic-field-induced splitting remains dominant in comparison with most of these mechanisms.

We would like to comment on the observability of the polariton zitterbewegung. Exciton–polaritons possess a finite lifetime which strongly bounds their propagation distance. Recent achievements in engineering of dielectric microcavities allow exciting polaritons characterized by lifetimes of the order of hundreds of picoseconds and propagating by the distances over 1 mm in the cavity plane [41–43]. However, these significant propagation distances are characteristic of photon-like polaritons. In our manuscript, we focus on exciton-like polaritons whose lifetime cannot exceed tens of picoseconds, realistically. To increase the distance of the macroscopic spreading of polaritons, we propose taking advantage of the stimulated scattering of excitons from the reservoir to the polariton mode [32]. The reservoir is conveniently created by a non-resonant optical pumping in the sub-threshold regime. This enables observing the polariton zitterbewegung over several periods of spatial oscillations. The presence of the exciton reservoir may trigger the energy relaxation of the polariton condensate [33]. This effect is expected to amplify the decay of the amplitude of oscillations with distance. Nevertheless, the zitterbewegung of polaritons remains observable. In the case of the resonant pumping alone, the spreading distance of the macroscopic polariton wave packet is strongly limited by the finite lifetime of propagating polaritons and hardly reaches several tens of micrometers. The color maps in figure 3 show that the system can be tuned to the conditions allowing for the observation of full-period oscillations even at such a small spreading distance. If the period of oscillations significantly exceeds the polariton propagation length, the effect of the zitterbewegung can be identified as the deviation of the wave packet trajectory from the linear ballistic trajectory. Still, these traces of the zitterbewegung are far less impressive than the long oscillating trajectories expected in the case of a combined resonant and non-resonant pumping (see figure 4).

The observation of the zitterbewegung of exciton–polaritons is of a fundamental importance as it gives the opportunity to observing the manifestation of a purely quantum mechanical phenomenon on a macroscopic scale. From the practical point of view, the magnetic control over the characteristics of the zitterbewegung as well as over the polarization properties of the polariton condensate allows considering the system as a polaritonic analog of the Datta-and-Das spin transistor [19, 44, 45]. The ability of separating in space of orthogonal linear diagonal polarizations allows one to use the proposed system as a polariton rectifier of the cw light.
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ORCID iDs

E S Sedov https://orcid.org/0000-0001-5684-151X

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