Exciton polaritons formed by strong coupling of excitons confined in quantum wells (QWs) and photons in high-quality optical microcavities (MCs) have been subject to active research for nearly three decades. The strong coupling results in the energy anticrossing of the exciton and photon modes, which generates two new normal modes known as the upper polariton (UP) and the lower polariton (LP) branches, as shown schematically in Fig. 1(a).

One of the widely used schemes for optical excitation of exciton polaritons involves off-resonant excitation by photons with energies well above the exciton energy in the QW material, which creates high-energy electron-hole pairs that subsequently lose their excess kinetic energy through interaction with lattice phonons to form bound pairs – the excitons. Once the kinetic energy of the excitons corresponds to momentum close to the light cone, strong coupling to the cavity photons allows for the formation of exciton polaritons. While the low-density, “thermal” population of the LP branch can be achieved through the phonon-assisted energy relaxation of high-energy carriers, it is not sufficient to observe macroscopic occupation of low-momentum states, i.e., bosonic condensation of exciton polaritons close to the ground state. The phonon relaxation becomes inefficient once the polariton reaches the inflection point of the dispersion and accumulate at high energy (momenta), in the so-called bottleneck region, where they form an incoherent reservoir of “hot” carriers with a large excitonic component. Further energy relaxation and macroscopic occupation of low-momentum states can only be achieved through stimulated bosonic scattering from the reservoir into the final states. This picture of the spontaneous condensation process has been widely accepted by the exciton-polariton community, and mean-field models reflecting this process have been successfully used to describe a vast variety of experiments. However, the question of whether the reservoir mostly consists of heavy, immobile excitons with negligible diffusion lengths or largely exciton-like, but highly mobile polaritons in the bottleneck region still has not been resolved. In some instances, a double-reservoir model is introduced to describe both populations, with one reservoir undergoing stimulated bosonic scattering into the condensate, and the other serving to replenish this gain medium and counteract its depletion. However, both of these reservoirs are usually considered to be immobile.

The light-induced excitonic reservoir is a key ingredient of a widely used optical trapping technique for exciton polaritons. Optical trapping of exciton polaritons achieved by structuring the pump using different methods such as amplitude masks, axicon lenses, Spatial Light Modulators (SLMs), and Digital Mirror Devices (DMD) has gained considerable attention, offering a platform for studying the exciton-polariton behaviour in arbitrary potentials. Polaritons interact due to the Coulomb repulsion of their excitonic component, and their interaction with the reservoir particles is repulsive. Therefore, the reservoir spatially localised in the pump region creates a local blueshift of the exciton-polariton energy and acts as a potential bar-
We use a high-quality GaAs/AlGaAs planar microcavity sample with 12 embedded GaAs quantum wells (QW) of 7 nm width sandwiched between distributed Bragg reflectors (DBR) of AlAs/AlGaAs with 32 (top) and 40 (bottom) layer pairs. The microcavity is operating in the strong coupling regime characterised by a vacuum Rabi splitting $\hbar \Omega = 15.9$ meV. The cavity photon lifetime is measured to be 135 ps resulting in polariton lifetimes of the order of $\sim 200$ ps. The sample is pumped nonresonantly with a continuous wave (CW) Ti:Sapphire laser at energies corresponding to high-energy reflectivity minima of the cavity structure. The sample is kept at $\sim 7$ K in a helium flow cryostat, and the photoluminescence (PL) is collected using a high numerical aperture (NA = 0.5) objective. A spectral edge-pass filter placed before the imaging CCD camera is used to filter out the laser light scattered from the sample. The overall experimental setup is similar to that used in Refs. 18 and 19.

We measure the PL by applying a circular filter in the far-field ($k$-space) image plane to collect only the PL of $k_\parallel \approx 0$ polaritons and to filter out the states with higher kinetic energy. This ensures that the spectral distribution of PL signal in real space measured along the monochromator slit, $E(x)$, effectively reflects the potential energy landscape for the polaritons. Hence, at low polariton densities, the spectrum represents the shape of the optically-induced potential generated by the reservoir\textsuperscript{19}. It should be noted that this filtering technique limits the spatial resolution of the setup due to the diffraction limit imposed by the $k$-space aperture, which is about 2 $\mu$m in our setup. Using a smaller aperture $k$-space filter would cause further broadening of the image in the conjugate real space plane, therefore we choose a moderate aperture size to avoid any deconvolution analysis of the real-space PL spectrum\textsuperscript{19}.

The experiments are performed on an area of the sample, where the photon-exciton detuning is positive, i.e. the cavity photon energy $E_c$ at zero momentum in the plane of the quantum well is larger than the exciton energy $E_X$: $\Delta(k = 0) = E_c(0) - E_X > 0$. This ensures that the exciton polaritons have a large excitonic fraction, which is quantified by the Hopfield coefficient $|X|^2 > 0.5\textsuperscript{10,12}$.

Following the studies of the exciton diffusion QWs\textsuperscript{26,29–31}, we model the spatial distribution of the incoherent reservoir as that of a classical gas of photoexcited carriers characterised by the effective diffusion length, $L_\text{eff}$. The spatio-temporal expansion of a locally excited distribution of carriers is then described by the classical diffusion equation:

$$\frac{\partial n(x,y,t)}{\partial t} = D\nabla^2 n(x,y,t) - \frac{n(x,y,t)}{\tau} + g(x,y,t), \quad (1)$$

where $n(x,y,t)$ is the density of the diffusing carriers at position $(x,y)$ and time $t$, $D$ represents the diffusion coefficient (diffusivity) for the carrier density, $\tau$ denotes the lifetime of the carriers, and $g(x,y,t)$ stands for the local rate of generation of the carriers.

An analytical solution for 2D stationary distribution of carriers created by a point source $g(x,y) = \delta(x-\xi)\delta(y-\eta)$ is given by\textsuperscript{32}:

$$G(x,y;\xi,\eta) = \frac{1}{2\pi D} K_0 \left( \frac{\sqrt{(x-\xi)^2 + (y-\eta)^2}}{\sqrt{D\tau}} \right), \quad (2)$$
where \( K_0 \) is the zeroth-order modified Bessel function of the second kind and the effective diffusion length is defined as \( L_{\text{eff}} = \sqrt{D\tau_h} \). In order to fit this distribution function to the spatial distribution of the carrier density in the experiment, one has to convolve Eq. (2) with the laser profile on the sample, which acts as an initial source of the carrier density. In our analysis, we assume that at low exciton-polariton densities the main contribution to the effective potential originates from the reservoir density, thus \( E_{\text{eff}}(x) = n_R(x) \). Therefore, the real-space PL spectrum filtered at \( k_{||} = 0 \) and corresponding to the effective potential allows us to extract the distribution of the reservoir density and its \( L_{\text{eff}} \).

We employ a 2D excitation scheme, where the optical pump is a focused Gaussian laser spot with a full width at half maximum (FWHM) of about 4 \( \mu \text{m} \). In this scheme, the laser creates a Gaussian-shaped potential “hill” (Fig. 1b). The spatially resolved PL spectra filtered at \( k_{||} = 0 \) are shown in Figs. 2(a-e) for various pump powers, \( P \), below the condensation threshold, \( P_{th} \). Intuitively, one expects that, as we pump stronger, more high-energy reservoir particles are injected into the system, which should result in an effective broadening of the reservoir distribution and increase of \( L_{\text{eff}} \). However, the analysis described above reveals two main effects of this excitation scheme. First, the effective diffusion length \( L_{\text{eff}} \) of the reservoir extracted from fitting the potential peak at the pump location decreases as a function of the increasing pumping power, as shown in Fig. 2(f). Secondly, the long tails of the PL distribution away from the pump spot are blueshifted from the lowest energy of the LP branch at \( k_{||} = 0 \) at a given spatial position, which indicates a non-negligible density of the reservoir and polariton-reservoir interaction up to the tens of microns away from the excitation spot (blueshift caused by polariton-polariton interactions is negligible at these low densities), see Fig. 3(c).

The apparent narrowing of the reservoir distribution at the location of the pump spot can be attributed to the local heating of the sample. The local increase of the sample temperature induces a narrowing of the semiconductor bandgap resulting in the lowering of the exciton energy and, additionally, temperature-induced change in the effective refractive index of the microcavity. This produces a local trapping potential for excitons and exciton polaritons, which reduces carrier mobility away from the pump spot. In previous studies the heating was linked to phonon-assisted thermal relaxation of the excitonic reservoir towards low-energy polaritons and therefore suggested as the mechanism for self-trapping of exciton polaritons at large densities, above the condensation threshold.

To confirm the presence of this heating effect in our low-density below-threshold regime, we have carried out acousto-optic modulator (AOM) duty cycle and frequency dependence measurements of the PL spectra at constant pumping powers to extract \( L_{\text{eff}} \). The AOM is usually used in cw experiments to chop the laser beam and reduce the heating of the sample. Typically, thermal effects relax much longer timescales (of the order of 10 – 100 \( \mu \text{s} \)) to that of the exciton-polariton dynamics (of the order of 1 – 100 ps). Increasing the duty cycle means that the duration of the pump pulse (comparable to the thermal relaxation timescale) increases, which directly affects the local temperature of the sample. First, we change the duty cycle from 5 to 99 \%, while keeping its frequency at 10 kHz, which corresponds to a period of \( T = 100 \mu \text{s} \) and the pulse duration 5 to 99 \( \mu \text{s} \), and measure the PL spectra at three different pump powers. The results are presented in Fig. 3(a). One observes that \( L_{\text{eff}} \) stays constant at the lowest pump power (\( P = 0.1 P_{th} \)) because at this power the local heating is negligible. However, at intermediate and large pump powers, the local heating and self-trapping of the reservoir carriers takes place and causes \( L_{\text{eff}} \) to decrease as with increasing duty cycle (pulse duration).

Next, we change the chopping frequency from 5 to 50 kHz and keep both the pumping power and the duty cycle fixed. Increasing the AOM frequency at a constant duty cycle decreases the pulse duration, which causes less heating at the pump spot thus a less pronounced self-trapping effect. As expected, Fig. 3(b) shows that the effective diffusion length increases with the increasing frequency, and saturates for AOM frequencies above 35 kHz (corresponding to about 30 \( \mu \text{s} \) period of 3 \( \mu \text{s} \) pulses). These results confirm our assertion that self-localisation of the carriers at the pump spot due to thermal effects is responsible for decreasing \( L_{\text{eff}} \) and hence narrowing potential barrier at larger pump powers.

The extended tails of the reservoir distribution away from the self-localised peak seen in Fig. 2(a-e) show the opposite tendency: this low-density distribution, i.e. the area where reservoir-induced blueshift of low-density polaritons is ob-
observed, broadens with increasing pump power. This results in a growing density of the reservoir at a fixed location away from the pump, as well as the increasing blueshift of the local polariton energy. This effect is illustrated in Fig. 3(c), where we plot the blueshift of the tails, similar to those shown in Figs. 2(a–e), at a fixed distance and pump power away from the pump region. These measurements are taken at a more excitonic detuning $\Delta = +10.5$ meV ($|X|^2 = 0.77$). As shown in Fig. 3(c), the local blueshift increases with increasing duty cycle, indicating that the local heating of the sample promotes reservoir transport over longer distances. The surprisingly long transport distances of the reservoir have been attributed to highly mobile bottleneck polaritons at high momenta in a previous study. Our observation, therefore, supports the suggestion that the reservoir responsible for the optically induced potential is composed of bottleneck polaritons with a very high excitonic component. On the other hand, the apparent growth of the reservoir background with temperature seen in Fig. 3(c) suggests that the reservoir transport could be assisted by phonons. Exciton drag by ballistically propagating non-equilibrium acoustic phonons (the so-called phonon wind effect) has been well studied in the literature (see, e.g., Ref. 26). In order to clarify the role of this mechanism in our experiment, further time-resolved measurements under a pulsed excitation will be required to establish the time-sequence of the self-focusing in the pump "hot spot" and long-range expansion of the reservoir. It would also be advantageous to design an experiment with controlled generation of the phonon wind, as suggested in Ref. 39.

To conclude, our analysis of the low-density polariton emission filtered at $k \approx 0$ reveals the details of the reservoir-induced potential landscape for exciton polaritons. Our results show a characteristic narrowing of the reservoir distribution at a pump-induced “hot spot”, which is consistent with the local bandgap renormalisation and formation of an effective self-localising potential for excitons. The low-density tails of the reservoir distribution extend several tens of micrometres away from the hot spot. This indicates that the transport of the reservoir particles, which have a large excitonic component, exceeds the transport length of bare excitons ($\sim 1 - 2$ $\mu$m) by an order of magnitude. These results are consistent with the recently reported extended transport of bottleneck polaritons. The physical reason for the enhanced transport, e.g. the role of the phonon wind, requires further careful investigation.

The observed behaviour of the reservoir is especially important when optical traps are employed to confine and manipulate exciton polaritons. The extended propagation distances and accumulation of the reservoir tens of micrometres away from the pump location affects local energy blueshifts of exciton polaritons and strongly contributes to the chemical potential of the polaritons above the bosonic condensation threshold. Our results indicate that the reservoir does not behave as an immobile density distribution strongly localised at the position of the pump, and this fact should be carefully taken into account when modelling exciton-polariton condensates in optical traps.

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FIG. 3. (a) Dependence of the $L_{\text{eff}}$ on AOM duty cycle for different pump powers $P = 0.1 P_{\text{th}}$ (circles), $P = 0.5 P_{\text{th}}$ (squares) and $P = 0.8 P_{\text{th}}$ (stars). AOM frequency is kept constant at 10 kHz. (b) Frequency dependence of $L_{\text{eff}}$ for $P = 0.8 P_{\text{th}}$ at AOM DC of 10%. Data taken at $\Delta = +5$ meV. (c) Blueshift of the tails $\approx 0$ emission at long distances away from the pump spot for changing AOM duty cycle range 30–50%, as indicated in the legend. Data taken at $\Delta = +10.5$ meV at low pumping power $P < P_{\text{th}}$.

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