

Quantum Droplets of Light in Semiconductor Microcavities

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Quantum droplets are dilute self-bound configurations of bosons that result from the balance between a mean-field attraction and a repulsion induced by quantum fluctuations. Such droplets have been successfully realized in cold atomic gases and represent a signature of their quantum nature. Here, we predict the existence of a similar droplet phase in a solid-state system, involving polaritons formed from the strong coupling between excitons (bound electron-hole pairs) and photons in a semiconductor microcavity. We consider a spin mixture of exciton-polaritons near a biexciton Feshbach resonance, which allows one to tune the interspecies interactions to be attractive and comparable in magnitude to the intraspecies repulsion. We find that self-bound quantum droplets are achievable for realistic parameters in atomically thin semiconductors, and that they can be detected via their excitation spectrum and spatial profile. This exotic phase could potentially lead to polariton condensation at lower thresholds and it opens an alternative avenue to achieve the long-sought quantum polaritonic regime.

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Classical liquids arise from the balance between two competing interatomic forces—a short-range repulsion and a long-range induced-dipole-dipole attraction [1]—and may form dense droplets held together by surface tension. Similar configurations also appear in the quantum realm in the form of dilute nanodroplets of liquid Helium [2], the prototypical strongly correlated quantum liquid [3]. Recently, a novel ultradilute self-bound quantum droplet was predicted to exist in ultracold mixtures of bosonic atoms with tunable interactions [4] (see also Ref. [5]). This arises due to a delicate competition between a net mean-field (MF) attraction—occurring when the interspecies attraction overcomes the intraspecies repulsion—and repulsive Lee-Huang-Yang (LHY) quantum fluctuations [6,7]. Remarkably, shortly after this was proposed, quantum droplets were experimentally realized in homonuclear [8,9] and heteronuclear [10] binary mixtures, as well as in dipolar condensates [11,12]. This has led to intense theoretical and experimental efforts [13], for instance on lowering the dimensionality [15–17], analyzing rotational [18,19] and collisional properties [20,21], considering new stabilizing mechanisms [22–25], and even different mixtures [26,27]. While quantum droplets have proved to be an extremely active research field, the focus has thus far been limited to ultracold atomic gases.

In this Letter, we predict the existence of self-bound quantum droplets in a solid-state system involving light and matter. Specifically, we consider exciton-polaritons, which are hybrid light-matter quasiparticles that emerge from the strong coupling between excitons (electron-hole bound states) and photons in a two-dimensional (2D) semiconductor microcavity [28–30]. Due to their bosonic nature and small effective mass, exciton-polaritons display remarkable collective coherent phenomena like Bose-Einstein condensation [28,31,32], superfluidity [33,34] and quantum vortices [35–37]. Over the last two decades, most experiments have been in the semiclassical regime where the collective effects can be accurately explained by a MF theory of non-linear classical waves [28,30]. However, recently, there has been a growing effort to access the regime of strongly correlated polaritons, which is necessary for a range of quantum photonic applications in scalable semiconductor systems [38–42]. Common proposals in this direction involve enhancing polariton-polariton repulsion (e.g., by exploiting dipolaritons [43–47] or Rydberg polaritons [48–50]), or strongly confining the cavity photon mode [51–53]. Here, we propose to instead *enhance* quantum fluctuations by effectively *weakening* the MF effects.

Key to our proposal is the ability to tune the interactions between polaritons, similarly to the case of ultracold atomic gases [8,9,11,12]. Specifically, we take advantage of the polariton Feshbach resonance [54–59], which arises when the energy of two opposite-spin polaritons matches that of a biexciton bound state [60–62]. As depicted in Fig. 1, this leads to a strong enhancement of interactions and a sudden

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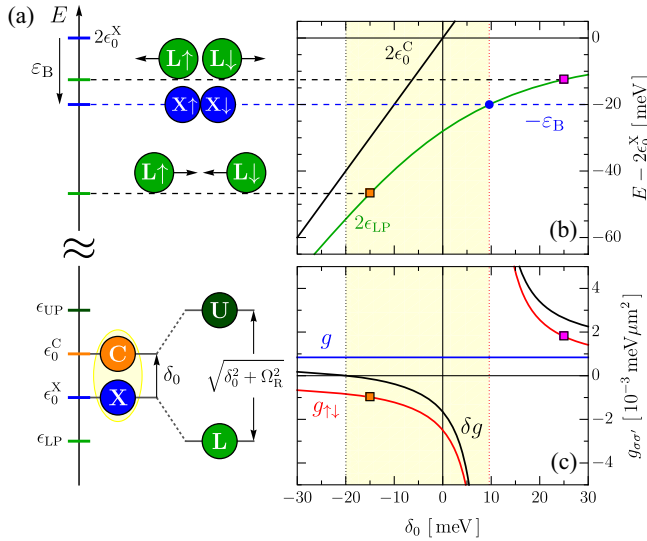


FIG. 1. (a) Schematic illustration of the energy scales for one (bottom) and two (top) particles. The hybridization of an exciton (X) and a photon (C) yields upper (U) and lower (L) polariton quasiparticles. Two $\uparrow$ and $\downarrow$ lower polaritons repel (attract) each other when their energy $2\epsilon_{\text{LP}}$ is above (below) the energy $-\epsilon_B$ of the biexciton ($X\uparrow-X\downarrow$). (b) Two-lower-polariton (green) and two-photon (black) energies as a function of detuning δ_0 . The blue point denotes the Feshbach resonance that occurs when $2\epsilon_{\text{LP}} = -\epsilon_B$. (c) Detuning dependence of the intra- (blue) and inter-species (red) exciton-exciton interaction coefficients. The magenta (orange) square is an example of repulsive (attractive) interspecies interactions. The black curve shows $\delta g = g + g_{\uparrow\downarrow}$, and the yellow region denotes the necessary condition $\delta g < 0$ for the appearance of quantum droplets. We use parameters inspired by recent experiments with MoSe_2 monolayers: $\Omega_R = 28$ meV, $\epsilon_X = 470$ meV [74], $\epsilon_B = 20$ meV [75], and exciton mass $m_X = 1.14m_0$ [76] (with m_0 the bare electron mass).

change in their sign, thus enabling us to enter a regime where quantum fluctuations dominate. The Feshbach resonance has recently been observed both in quantum well (QW) semiconductors [63–65] and in transition metal dichalcogenide (TMD) monolayers [66,67], and it has also received significant attention in the context of many-body effects such as Bose polarons [66,68–70], photon dimers [71], and novel collective paired phases [72,73].

Using a Bogoliubov approximation combined with a bosonic pairing approach [77–79], we demonstrate that the quantum droplet phase is robust over a wide range of experimentally accessible parameter regimes and can exist whenever mean-field attraction is present. We obtain several experimental signatures of this novel quantum polariton phase such as the excitation spectrum, the saturation density, and the spatial profile. In particular, we predict that the droplet will occur for parameters that are within reach of current experiments on semiconductor microcavities.

Model—We consider a 2D spin mixture of exciton-polaritons described by the Hamiltonian $\hat{H} = \hat{H}_0 + \hat{V}$

written in the photon-exciton basis (setting $\hbar$ and the system area to 1):

$$\hat{H}_0 = \sum_{\mathbf{k}\sigma} \begin{pmatrix} \hat{c}_{\mathbf{k}\sigma}^\dagger & \hat{x}_{\mathbf{k}\sigma}^\dagger \end{pmatrix} \begin{pmatrix} \epsilon_{\mathbf{k}}^C - \mu_\sigma & \Omega_R/2 \\ \Omega_R/2 & \epsilon_{\mathbf{k}}^X - \mu_\sigma \end{pmatrix} \begin{pmatrix} \hat{c}_{\mathbf{k}\sigma} \\ \hat{x}_{\mathbf{k}\sigma} \end{pmatrix}, \quad (1a)$$

$$\hat{V} = \sum_{\sigma\sigma'} \sum_{\mathbf{k}\mathbf{k}'\mathbf{q}} \frac{v_{\sigma\sigma'}}{2} \hat{x}_{\mathbf{k}+\mathbf{q}\sigma'}^\dagger \hat{x}_{\mathbf{k}'-\mathbf{q}\sigma'}^\dagger \hat{x}_{\mathbf{k}'\sigma'} \hat{x}_{\mathbf{k}\sigma}. \quad (1b)$$

Here, $\hat{x}_{\mathbf{k}\sigma}^\dagger$ ($\hat{c}_{\mathbf{k}\sigma}^\dagger$) creates an exciton (cavity photon) with in-plane momentum $\mathbf{k}$, spin (polarization) $\sigma = \{\uparrow, \downarrow\}$ and dispersions $\epsilon_{\mathbf{k}}^X = k^2/2m_X$ and $\epsilon_{\mathbf{k}}^C = k^2/2m_C + \delta_0$, respectively, where the typical photon to exciton mass ratio is $m_C/m_X = 10^{-5}$. We measure energies with respect to the $\mathbf{k} = 0$ exciton, with δ_0 the photon-exciton detuning, and μ_σ the chemical potentials for each spin. Equation (1a) treats the excitons as structureless bosons which is valid as long as the light-matter coupling is much smaller than the exciton binding energy, i.e., $\Omega_R \ll \epsilon_X$, as is the case in microcavities with either monolayer TMDs or a few III-V quantum wells.

The noninteracting Hamiltonian $\hat{H}_0$ is diagonalized by the transformations $\hat{L}_{\mathbf{k}\sigma} = C_{\mathbf{k}} \hat{c}_{\mathbf{k}\sigma} + X_{\mathbf{k}} \hat{x}_{\mathbf{k}\sigma}$, $\hat{U}_{\mathbf{k}\sigma} = X_{\mathbf{k}} \hat{c}_{\mathbf{k}\sigma} - C_{\mathbf{k}} \hat{x}_{\mathbf{k}\sigma}$, where $\hat{L}_{\mathbf{k}\sigma}$ and $\hat{U}_{\mathbf{k}\sigma}$ are the lower (LP) and upper (UP) polariton quasiparticles with dispersions $E_{\text{UP/LP}}(\mathbf{k}) = \frac{1}{2}(\epsilon_{\mathbf{k}}^C + \epsilon_{\mathbf{k}}^X \pm \sqrt{(\epsilon_{\mathbf{k}}^C - \epsilon_{\mathbf{k}}^X)^2 + \Omega_R^2})$, while $X_{\mathbf{k}}/C_{\mathbf{k}}$ are the Hopfield coefficients which quantify the exciton/photon fraction inside each quasiparticle [their explicit expressions can be found in Supplemental Material (SM) [80]]. In the following, we denote the polariton energies at zero momentum as $\epsilon_{\text{UP/LP}}$.

Polaritons inherit interactions from their exciton component, which have underlying spin-dependent interaction strengths $v_{\sigma\sigma'}$, as in Eq. (1b). Since exciton interactions are short-ranged, they can be described by an effective interaction strength $g_{\sigma\sigma'}$, which is obtained from the Born series expansion [80]

$$g_{\sigma\sigma'} = v_{\sigma\sigma'} \left(1 + v_{\sigma\sigma'} \sum_{\mathbf{k}} \frac{1}{2\epsilon_{\text{LP}} - 2\epsilon_{\mathbf{k}}^X} + \dots \right), \quad (2)$$

where the collision energy in the scattering is set by the two interacting lower polaritons [59]. For the repulsive intra-species interactions $g_{\uparrow\uparrow} = g_{\downarrow\downarrow} \equiv g$, it is sufficient to expand up to second order and use the standard Born approximation [84,85] for the underlying interactions $v_{\uparrow\uparrow} = v_{\downarrow\downarrow} \equiv v$. We take $g = 4\pi/m_X$ since this approximately corresponds to the values for both TMDs and QW semiconductors (within a $\sim 10\%$ uncertainty [86]).

On the other hand, the underlying interspecies interactions are attractive, i.e., $v_{\uparrow\downarrow} = v_{\downarrow\uparrow} < 0$, and they lead to a $\uparrow\downarrow$ biexciton bound state which requires the infinite Born series in Eq. (2). Importantly, in the regime $\Omega_R \ll \epsilon_X$ the

effective interaction $g_{\uparrow\downarrow}$ takes a universal low-energy form in the vicinity of the biexciton bound state with binding energy ϵ_B [59]

$$\frac{1}{g_{\uparrow\downarrow}} = \frac{1}{v_{\uparrow\downarrow}} + \sum_{\mathbf{k}} \frac{1}{2\epsilon_{\mathbf{k}}^X - 2\epsilon_{\text{LP}}} = \frac{m_X}{4\pi} \ln\left(\frac{\epsilon_B}{2|\epsilon_{\text{LP}}|}\right). \quad (3)$$

In contrast to 2D ultracold gases, $g_{\uparrow\downarrow}$ depends on the light-matter coupling instead of the system density.

As shown in Fig. 1, the interspecies interaction can be tuned from attractive ($g_{\uparrow\downarrow} < 0$) to repulsive ($g_{\uparrow\downarrow} > 0$) by varying the photon detuning δ_0 across the biexciton Feshbach resonance $\epsilon_{\text{LP}} = -\epsilon_B/2$ [87]. For concreteness, we use typical experimental parameters for TMD monolayers. While a collapse instability of the polariton condensate is predicted at the MF level [88] when $\delta g \equiv g + g_{\uparrow\downarrow} < 0$ [yellow region in Figs. 1(b) and 1(c)], we show here that beyond-MF quantum fluctuations can stabilize the system into a self-bound droplet phase. Note that since the Born approximation represents an upper bound on g [89], in reality we expect the yellow region to extend to larger negative detunings. The precise value of the lower critical detuning can also depend on the finite range of the interactions.

Bogoliubov theory with bosonic pairing—We now turn to the many-body system and consider the thermodynamic potential $\Omega = \langle \hat{H} \rangle_0$, where $\langle \dots \rangle_0$ denotes an expectation value over the Bose-Einstein condensate (BEC) of LPs. To capture the biexciton Feshbach resonance, we introduce the pairing field $\Phi = -v_{\uparrow\downarrow} \sum_{\mathbf{p}} \langle \hat{x}_{\mathbf{p}\uparrow} \hat{x}_{-\mathbf{p}\downarrow} \rangle$ and we perform a MF decoupling of the singlet sector of the interaction term. Thus, Eq. (1b) becomes

$$\hat{V} = \sum_{\mathbf{k}\mathbf{k}'\mathbf{q},\sigma} \frac{v}{2} \hat{x}_{\mathbf{k}+\mathbf{q}\sigma}^\dagger \hat{x}_{\mathbf{k}'-\mathbf{q}\sigma}^\dagger \hat{x}_{\mathbf{k}'\sigma} \hat{x}_{\mathbf{k}\sigma} - \frac{|\Phi|^2}{v_{\uparrow\downarrow}} - \sum_{\mathbf{k}} (\Phi \hat{x}_{\mathbf{k}\uparrow}^\dagger \hat{x}_{-\mathbf{k}\downarrow}^\dagger + \Phi^* \hat{x}_{-\mathbf{k}\downarrow} \hat{x}_{\mathbf{k}\uparrow}). \quad (4)$$

A similar bosonic pairing approach was first considered in the context of attractive single-component Bose liquids [77], and has already been used to model quantum droplets in 3D Bose mixtures [78,79].

The macroscopic condensation of polaritons into the zero-momentum state can be described by a zero-temperature semiclassical approximation of the exciton and photon operators, i.e., $\hat{x}_{0\sigma} \sim \alpha_\sigma$ and $\hat{c}_{0\sigma} \sim \lambda_\sigma$. Since the system Hamiltonian $\hat{H}$ is endowed with a $U(1) \times U(1)$ symmetry, i.e., it is invariant under a global transformation $\alpha_\sigma(\lambda_\sigma) \mapsto e^{i\theta_\sigma} \alpha_\sigma(\lambda_\sigma)$ and $\Phi \mapsto e^{i(\theta_\uparrow + \theta_\downarrow)} \Phi$, we choose $\alpha_\sigma, \lambda_\sigma, \Phi \in \mathbb{R}$ without any loss of generality. While a general derivation can be found in the SM [80], here we focus on the experimentally relevant case of a symmetric spin mixture ($\alpha_\sigma = \alpha/\sqrt{2}$, $\lambda_\sigma = \lambda/\sqrt{2}$ and $\mu_\sigma = \mu$), where

the thermodynamic potential at the mean-field level reads,

$$\Omega_{\text{MF}} = -\frac{\Phi^2}{g_{\uparrow\downarrow}} - (\Phi + \mu)\alpha^2 + (\delta_0 - \mu)\lambda^2 + \Omega_R \lambda \alpha + \frac{g}{4}\alpha^4, \quad (5)$$

in terms of g and $g_{\uparrow\downarrow}$. Since photons are noninteracting, the photon densities enter Ω only through its MF part, so the stationary condition $\partial_\lambda \Omega_{\text{MF}} = 0$ relates them to the exciton densities as $\lambda = -\Omega_R \alpha / [2(\delta_0 - \mu)]$. This relation, which in the noninteracting case determines the total occupation of the LP branch [80], can be inserted into Eq. (5) to give

$$\Omega_{\text{MF}}(\alpha, \Phi, \mu) = -\frac{\Phi^2}{g_{\uparrow\downarrow}} - \left(\Phi + \mu + \frac{\Omega_R^2}{4(\delta_0 - \mu)} \right) n_0 + \frac{g}{4} n_0^2, \quad (6)$$

where $n_0 = \alpha^2$ is the exciton density of the polariton condensate. Ω_{MF} explicitly depends on the chemical potential, whose MF value μ_{MF} is obtained by minimizing Ω_{MF} with respect to α [80].

The next contribution is the LHY term coming from quantum fluctuations around the condensate. This is evaluated at $\mu = \mu_{\text{MF}}$ and reads,

$$\Omega_{\text{LHY}}(\alpha, \Phi) = \frac{1}{2} \sum_{\mathbf{k} \neq 0} \left[E_{+, \mathbf{k}}^a + E_{-, \mathbf{k}}^a + E_{+, \mathbf{k}}^b + E_{-, \mathbf{k}}^b - 2(A_{\mathbf{k}} + B_{\mathbf{k}}) + \frac{C_+^2 + C_-^2}{2\epsilon_{\mathbf{k}}^X + 2|\epsilon_{\text{LP}}|} \right]. \quad (7)$$

Here, $E_{\pm, \mathbf{k}}^{a/b}$ are the four Bogoliubov branches of the system and we have introduced $A_{\mathbf{k}} = \epsilon_{\mathbf{k}}^C - \mu_{\text{MF}}$, $B_{\mathbf{k}} = \epsilon_{\mathbf{k}}^X - \mu_{\text{MF}} + g n_0$, and $C_{\pm} = g n_0 / 2 \pm \Phi$. The spectrum of excitations is obtained from the diagonalization of the Bogoliubov Hamiltonian, and analytical expressions of the four branches are found for a symmetric mixture. In the limit of small photon mass $m_C \ll m_X$, the Bogoliubov spectrum simplifies and the momentum integral in Eq. (7) can be computed analytically [80]. Notice that the validity of the Bogoliubov theory improves by moving further from the biexciton resonance, i.e., for negative values of detuning (cf. Fig. 1).

Droplet phase—The existence of self-bound quantum droplets relies on the coexistence between vacuum (zero density) and a finite-density phase at zero pressure P . Since the total thermodynamic potential $\Omega = \Omega_{\text{MF}} + \Omega_{\text{LHY}} = -P$ depends on the exciton density n_0 (via α), the pairing field Φ and the chemical potential μ , the quantum droplet state QD $\equiv (n_0, \Phi_D, \mu_D)$ is thus the simultaneous solution of the three conditions

$$(i) \Omega|_{\text{QD}} = 0, \quad (ii) \partial_\alpha \Omega|_{\text{QD}} = 0, \quad (iii) \partial_\Phi \Omega|_{\text{QD}} = 0 \quad (8)$$

at nonzero n_0 . To determine the existence of the droplet phase, we plot in Fig. 2(a)–2(d) the thermodynamic

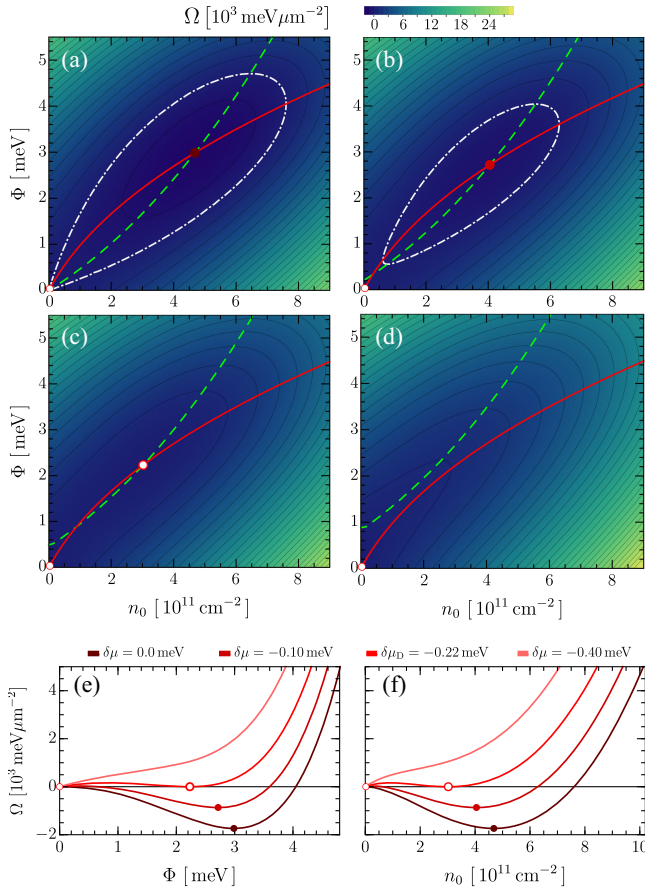


FIG. 2. (a)–(d) Thermodynamic potential as a function of exciton density and pairing field for 4 values of the chemical potential $\delta\mu = \mu - \epsilon_{LP}$: 0 (a), -0.10 meV (b), -0.22 meV (c), -0.40 meV (d). The white dot-dashed contour denotes the zero pressure line, i.e., the zeros of Ω , while the green dashed and red solid curves correspond to the conditions $\partial_a \Omega = 0$ and $\partial_\Phi \Omega = 0$, respectively. The filled dots in panels (a) and (b) are the stationary points satisfying both conditions. The empty dot in panel (c) indicates the quantum droplet which appears as a first order transition, signaling phase coexistence with the vacuum (white-red point). Thermodynamic potential as a function of pairing field (e) and exciton density (f) for the same values of $\delta\mu$, obtained by evaluating Ω along the red solid curve of the corresponding panels (a)–(d). The detuning is fixed at $\delta_0 = 0$, while all the other system parameters are the same as in Fig. 1.

potential in the n_0 - Φ plane at fixed chemical potential for the specific case of zero detuning—these results are extended to the full range of the quantum droplet phase ($\delta g < 0$) in Fig. 3. Here, we parametrize μ as a deviation from its non-interacting value, i.e., $\mu = \epsilon_{LP} + \delta\mu$.

Figure 2(a) shows the case $\delta\mu = 0$ where the intersections between the green ($\partial_a \Omega = 0$) and red ($\partial_\Phi \Omega = 0$) curves yield two stationary points: the vacuum, and a global minimum at finite density and positive pressure. This is apparent in Figs. 2(e) and 2(f), which show the behavior of Ω along the red curve as a function of the pairing field and the exciton density. Taking instead $\delta\mu = -0.10$ meV in

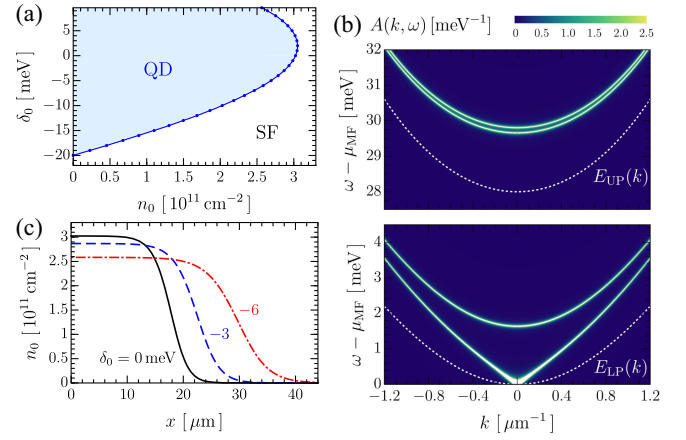


FIG. 3. (a) Phase diagram of detuning versus exciton density in the polariton condensate. The blue solid line bounds the region (QD) where quantum droplets can form, and it corresponds to the saturation density of the quantum droplet in phase coexistence with the vacuum. Outside this region, the system is in a superfluid miscible phase (SF). Quantum droplets only occur in the detuning range $-19.97 \text{ meV} < \delta_0 < 9.6 \text{ meV}$, where $\delta g \leq 0$ for our parameters. (b) Absorption spectrum, featuring the four Bogoliubov branches $E_{-k}^{a/b}$ (bottom panel) and $E_{+k}^{a/b}$ (top panel), which includes a linear gapless mode. The LP and UP dispersions are plotted as white dashed lines. The cavity linewidth has been fixed to $\Gamma = 0.025 \text{ meV}$. (c) Droplet density profiles, obtained from a Gross-Pitaevskii (GP)-like approach (see text), for $\delta_0 = 0.0$ (black, solid line), -3.0 (blue dashed) and -6.0 meV (red dash-dotted). System parameters are as in Fig. 1.

Fig. 2(b), the line of zeros shrinks and the vacuum becomes an isolated zero. As a result, the corresponding curves in Figs. 2(e) and 2(f) display a saddle point and a global minimum at negative and positive pressures, respectively.

For the critical value $\delta\mu_D \simeq -0.22 \text{ meV}$ in Fig. 2(c), the region of positive pressure disappears as the line of zeros contracts to a single point with $n_{0D} = 3.0 \times 10^{11} \text{ cm}^{-2}$ and $\Phi_D = 2.2 \text{ meV}$. This corresponds to phase separation between the vacuum and a polariton condensate at a particular density—the so-called saturation density—thus yielding a quantum droplet of exciton-polaritons. On the other hand, in Fig. 2(d) where $\delta\mu < \delta\mu_D$, the vacuum is the only stable phase.

As depicted in Fig. 3(a), we find that the quantum droplet exists with a well-defined saturation density across the full range of cavity detunings where $\delta g \leq 0$ [Fig. 1(c)]. Above this density, the system forms the usual superfluid phase without any instability. However, if polaritons are pumped such that the exciton density n_0 lies within the quantum-droplet region of the phase diagram in Fig. 3(a), then the system increases its density to form a quantum droplet surrounded by vacuum. This suggests the tantalizing possibility of achieving polariton lasing at a lower pump power.

At $\delta_0 \simeq -20 \text{ meV}$, the saturation density hits zero and only the miscible superfluid phase exists, while above the

Feshbach resonance at $\delta_0 > 9.6$ meV, the ground state is unstable towards forming biexcitons. Note that the non-monotonic behavior of the phase boundary in Fig. 3(a) can be ascribed to the presence of the pairing field since a similar behavior was observed in previous studies of pairing in quantum mixtures [72,73,90]. Phase diagrams as a function of either the pairing field or the chemical potential can be found in [80].

Excitation spectrum—Another key observable of the droplet phase is the excitation spectrum which can be accessed experimentally in polariton systems [91–95]. We plot in Fig. 3(b) the excitation absorption spectrum, which is proportional to the photon spectral function $A(\mathbf{k}, \omega) = -(1/\pi) \text{Im} [G_C(\mathbf{k}, \omega + i\Gamma)]$ [96], where the photon Green’s function is defined in [80] and Γ is the cavity photon linewidth. We observe that the lowest Bogoliubov branch is linear and gapless, i.e., $E_{-\mathbf{k}=0}^a = 0$, as a consequence of the $U(1)$ symmetry breaking of the system. The pairing field in our theory leads to a finite energy gap in the next branch $E_{-\mathbf{k}}^b$, which one expects will be removed once further fluctuations are included. The spectrum in Fig. 3(b) also shows two additional branches which lie higher in energy, as they result from the splitting of the UP branch. A clear experimental signature of the polariton quantum droplet phase is the behavior of the speed of sound of the gapless mode, since it remains constant for particle numbers below the saturation density but increases in the miscible superfluid phase.

Droplet profiles—Finally, we estimate the size and spatial profile of polariton quantum droplets using a Gross-Pitaevskii (GP) approach within a slab-geometry configuration [97,98]. Here we assume a sizable condensate with a uniform density along the y direction and a liquid-vacuum interface along the x axis (see Ref. [80] for details), thus yielding an upper bound for the droplet size. As a result, we find that the spatial profile of the exciton density $n_0(x)$ is governed by the equation

$$\frac{\hbar^2}{8M_*} \left[-2 \frac{n_0''(x)}{n_0(x)} + \left(\frac{n_0'(x)}{n_0(x)} \right)^2 \right] + \frac{\partial \Omega}{\partial n_0} [n_0(x)] = 0, \quad (9)$$

where the light-matter coupling is accounted for by the effective mass $M_*^{-1} = m_X^{-1} + \Omega_R^2/[4m_C(\delta_0 - \mu)^2]$, which is linked to the LP effective mass via $M_* \simeq X_0^2 M_{\text{LP}}$ [30]. Here we have eliminated the dependence on the pairing field using the stationary condition (8) (iii). Note that the interplay between the inter- and intraspecies interactions is encoded within the density derivative of the thermodynamic potential which thus results in an effective “trapping potential” [80]. Figure 3(c) shows how the density profiles have a flat top, and that the saturation density and droplet radius decreases and increases, respectively, as $|\delta g|$ becomes smaller (i.e., considering more negative cavity detuning). Overall, we find droplet radii of order ~ 10 μm , which is compatible with ongoing experiments.

Conclusions—In this Letter, we have shown that exciton-polaritons can host quantum droplets, the ultimate liquid-like behavior of a quantum fluid of light. Our findings complement previous studies of liquid light in the context of nonlinear optics with competing nonlinearities [99–102], as well as quasiequilibrium photon fluids in plasma and semiconductor environments [103]. The advantage of polariton droplets over their cold-atom counterparts is that they do not suffer from three-body recombination, due to the lack of deep two-body bound states [104], and there is in principle a greater range of particle numbers that can be accessed. Furthermore, the Bogoliubov excitation spectrum is more readily probed in polariton condensates, thus providing a deeper insight into the broken symmetries and associated Goldstone modes of a Bose mixture. Specifically, there is currently a lack of consensus between different perturbative methods [4,78] about the dispersions of the spin and density modes. The driven-dissipative nature of the polariton system is not expected to change our results qualitatively for long polariton lifetimes [72,92,105] and these effects can also be straightforwardly incorporated in our GP analysis, e.g., through a nonresonant pumping scheme [106].

Our results suggest that quantum droplets of light are achievable for realistic ranges of parameters in current TMD experiments. Key signs of the droplet phase include an ultra-low-threshold for condensation, the absence of a blueshift with varying pump power, an unchanging Bogoliubov spectrum, and a condensate size of ~ 10 μm . The observation of quantum droplets would provide an unmistakable signature of the elusive quantum nature of exciton-polariton quasiparticles, thus opening promising avenues into the long-standing challenge of quantum polaritonics [41,42].

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Data availability—The data that support the findings of this article are openly available [107].

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