Quasi-condensation of bilayer excitons in a periodic potential

Camille Lagoin$^1$, Stephan Suffit$^2$, Kenneth West$^3$, Kirk Baldwin$^3$,
Loren Pfeiffer$^3$, Markus Holzmann$^4$ and François Dubin$^1$

$^1$ Institut des Nanosciences de Paris, CNRS and Sorbonne Université, 4 pl. Jussieu, 75005 Paris, France
$^2$ Laboratoire de Matériaux et Phénomènes Quantiques, Université Paris Diderot, 75013 Paris
$^3$ PRISM, Princeton Institute for the Science and Technology of Materials, Princeton University, Princeton, NJ 08540 and
$^4$ Univ. Grenoble Alpes, CNRS, LPMMC, 38000 Grenoble, France

We study two-dimensional excitons confined in a lattice potential, for high fillings of the lattice sites. We show that a quasi-condensate is possibly formed for small values of the lattice depth, but for larger ones the critical phase-space density for quasi-condensation rapidly exceeds our experimental reach, due to the increase of the excitons effective mass. On the other hand, in the regime of a deep lattice potential where excitons are strongly localised at the lattice sites, we show that an array of phase-independent quasi-condensates, different from a Mott insulating phase, is realised.

The phase diagram of bosonic particles exploring a lattice potential is commonly described in terms of two essential parameters, the inter-particle interaction strength $U$ giving the increase of energy when two particles occupy the same lattice site, and the particle tunnelling strength $J$, determining the hopping of a particle to a neighbouring site [1, 2]. Within the minimal Bose-Hubbard model, two antagonist regimes emerge at zero temperature. When tunnelling is dominant, a superfluid showing phase ordering over spatially distinct sites can be reached, whereas above a critical ratio $U/J$ tunnelling is suppressed and bosons remain fixed to their lattice site. This results in an incompressible phase, a so-called Mott insulator, where phase coherence extending over various lattice sites is absent. Typically, varying the lattice depth or the lattice period $U$ and $J$ are modified and the phase diagram is accurately explored [1, 2].

Seminal experiments conducted with ultra-cold atomic gases confined in optical lattices have provided model studies of the Bose-Hubbard model, initially in three dimensions, and more recently, for effectively two and one dimensional systems [2]. In the solid-state, large efforts have also been dedicated to the physics of the two-dimensional Bose-Hubbard Hamiltonian, in particular for quantum simulation perspectives. To this aim, promising candidates include long-lived dipolar excitons of semiconductor bilayers. Such excitons are marked by a spatial separation imposed between Coulomb-bound electrons and holes. They are possibly engineered in coupled GaAs quantum wells [3] or in van der Waals assemblies of transition metal dichalcogenides monolayers [4]. In the former case, excitons may be subject to an artificial lattice potential engineered by electrostatic gates [5, 6] while in the latter case they naturally explore a periodic moire potential [7–9]. For both systems, the Bose-Hubbard model is expected to provide an accurate low-energy description.

To explore the collective phases accessible to dipolar excitons, GaAs double quantum wells provide an experimental toy model system, since excitons are then possibly manipulated in a regime of homogeneous broadening [10], and spatially confined in on-demand potential landscapes [11–16]. Using these degrees of freedom, we have recently mapped out the quasi-condensation crossover of bilayer excitons in box-like trapping potentials [10, 17, 18], determining the excitons equation of state and density fluctuations, and correlated these to the degree of spatial and temporal coherence at sub-Kelvin temperatures.

Here, we report experiments characterising the quasi-condensation of GaAs bilayer excitons in a microscopic lattice potential. We show that a quasi-condensate is annihilated above a threshold lattice depth around a fraction of the chemical potential. We attribute the suppression of coherence to the renormalisation of the excitons effective mass by the lattice potential which rapidly increases the critical phase space density of quasi-condensation. Above a threshold lattice depth, cold excitons enter then a normal phase, but without yet being spatially localised at the lattice sites. The localised regime is in fact obtained for lattice depths large compared to the chemical potential. In this regime, we show that an array of phase incoherent quasi-condensates localized at each site is possibly formed at large filling factors. We underline that this realisation differs from a Mott insulator.

In our studies, dipolar excitons are created in a square lattice with spatial period $L=3 \mu m$, at a bath temperature $T_b$ set to 340 mK (see Ref. [5] for a complete description of our device). We focus onto excitonic properties at a fixed average density $n \sim 2 \times 10^{10}$ cm$^{-2}$ such that the phase-space density $D$ is around 12 while the excitons chemical potential $\mu$ is about 0.8 meV. To quantify the role of the lattice depth on the excitons quasi-condensation, we continuously vary the height of the potential barrier separating lattice sites $V_0$, from 0 to 3 meV, by adjusting the potential difference applied between a pair of interdigitated electrodes deposited at the surface of the field-effect device embedding the double GaAs quantum well [5].

Figure 1.d illustrates the phases potentially accessible in our experiments: for small $U/J$, realized in the limit of vanishing lattice depth, a quasi-condensate phase marked by quasi-long range order in the spatial and temporal coherence is expected [1, 19]. On the other hand, for a deep lattice potential, i.e. for large
We note that initially it decays exponentially with a characteristic time constant of 4 µs. The dashed line shows an algebraic time decay function of the time delay |τ|∼ τc with |τc|∼ 4 is a dimensionless parameter quantifying dipolar interactions between excitons [10]. For our experimental conditions we deduce that U∼ kBTb, so that Mott phases are not energetically protected and thermal occupation of higher bands of the periodic lattice cannot be neglected.

We now turn to discuss our experimental results in these different regimes. Let us start with the quasicondensate regime at vanishing U/J, by setting V0 to 0. In this situation, we studied the degree of spatial and temporal coherence of the photoluminescence radiated by our device using a Mach-Zehnder interferometer where the photoluminescence field ψ is recombined with itself, after a variable time delay τ, or a spatial displacement r are introduced. Measuring the interference contrast at the interferometer output, we directly deduced the amplitude of the first-order correlation function |g(1)(r, τ)|∼ |⟨ψ∗(r, t)ψ(r + r, t + τ)⟩|, where ⟨·⟩t denotes the average over t, i.e. over the number of realisations accumulated to produce one interferogram, as well as the average over the positions r, where the fringe contrast is evaluated, typically of the order of 10 µm to compute the contrast from 3 fringes (Fig.2.a-b).

FIG. 1: (a-c) Real images of the photoluminescence for V0=0 (a), 1.5 meV (b) and 3 meV (c). Scales are given in µm. (d) Schematic phase diagram for the two-dimensional Bose-Hubbard model, as a function of the on-site interaction U, the tunnelling strength J and the temperature T. At non-zero temperature superfluid and Mott insulating phases, SF and MI respectively, are separated by a normal phase where physical properties are classical. Our experiments, carried out at a fixed average density, explore the parameter space along the dashed curved line. (e) At the crossover between the SF and normal phases non-condensed excitons (blue) are mostly located around the lattice barrier. This effectively increases the lattice depth (solid line) compared to the native one (dashed line). Quasi-condensed excitons are depicted in red.

FIG. 2: (a) Interference pattern measured for V0=0 and for (r=0, τ=4.2 ps), together with the profile evaluated at the center of the image (b) leading to |g(1)|=28 %.
(c) Decay of the first-order time coherence |g(1)| as a function of the time delay τ. Experimental data are depicted by the solid blue points, while the red line shows an exponential decay with a characteristic time constant of 4 ps. The dashed line shows an algebraic time decay τ−n with n=0.25. The inset presents the same measurements in linear scale. (d) Spatial interference of the photoluminescence measured for (|r|=2 µm, τ=0), together with the interference profile evaluated at the center of the image (e). Thus we deduce an interference contrast |g(1)|=23 %.

In (a-b) the interference period is set to 3 µm and 4.5 µm in (d-e). Experiments have all been performed at a bath temperature Tb=340 mK.

Figure 2.c shows the variation of the photoluminescence first order time correlation function |g(1)|, We note that initially it decays exponentially with a time constant τc∼ 4 ps, before a slower decay is found.
for $4 \lesssim \tau \lesssim 7$ ps. Very recently we have reported such behaviour and shown that the initial exponential decay reveals the contribution of non-condensed excitons through the inelastic two-body collisional rate. Indeed, the exciton-photon coupling is linear so that the coherence of optically bright excitons is imprinted in the photoluminescence [20]. Thus, we find here an excitons collisional rate, $1/\tau_c$, in good agreement with our previous studies [10]. On the other hand, the slowly decaying part for $\tau \gtrsim \tau_c$ marks the contribution of quasi-condensed excitons, which exhibit an algebraically decaying time coherence $\tau^{-\eta}$ [19]. Importantly, here we find that $\eta \sim 0.25$ which is compatible with the value predicted at criticality by the Berezinskii-Kosterlitz-Thouless theory [21, 22]. Let us then note that our studies are realised for an exciton phase-space density $D \sim 12$, very close to Ref. [10] where a very similar exponent was deduced.

We further studied the degree of spatial coherence by setting our interferometer such that $|\langle r \rangle| \sim 2 \mu m$ and $\tau=0$. This spatial separation is over an order of magnitude beyond the classical limit set by the thermal de Broglie wavelength, and two times larger than our optical resolution. Figure 2.d reveals that we observe then interference fringes signaling the buildup of quasi-long-range spatial coherence. In these experiments the interference contrast amounts to 23%, compared to 75% at the spatial auto-correlation. From this difference we deduce that $1/3$ of the optically bright excitons population contribute to the quasi-condensate [23]. In previous works, we have shown that in GaAs a quasi-condensate also includes optically dark excitons, that do not contribute to the photoluminescence but constitute around $2/3$ of quasi-condensed excitons [24]. Then, we deduce that about 75% of the total exciton population undergoes quasi-condensation in our experiments. Importantly this magnitude agrees with the one we previously deduced in electrostatic traps for similar densities [17], in good agreement with path integral Monte-Carlo calculations [25]. Also, we wish to stress that the amplitude of $|\langle r^{(1)} \rangle|$ is of the same order for $|\langle r \rangle| \sim 2 \mu m$, $\tau=0$ and $\tau=0$, $\tau \sim 6$ ps. This matching is expected since in both cases only the fraction of quasi-condensed bright excitons contributes to the interference signal.

We now turn to the opposite regime where a deep lattice potential is imprinted, $V_0=3$ meV $\gg \mu$. In this situation Fig.1.c shows that excitons are strongly localised in the lattice sites which have a characteristic spatial extension of about $1 \mu m$. This results in a modulation of the photoluminescence intensity, of around 20% along the horizontal axis of our device, with a 3 $\mu m$ period [5]. To study spatially extended time coherence in the photoluminescence, we studied $|\langle r^{(1)}(0, \tau) \rangle|$ fixing the interference period to 4.5 $\mu m$, as in the measurements shown in Fig. 2.d-e. Thus, the interference period is clearly distinguished from the one of the spatial modulation of the photoluminescence intensity due to the lattice potential. Figure 3.a shows the variation of the interference contrast with $\tau$, which, unlike in Fig.2.c, is reduced to an exponential decay with a characteristic time of around 3 ps. We then note that for $V_0=3$ meV a slightly higher rate of inelastic collisions is deduced compared to $V_0=0$. We attribute this difference as the manifestation of an increased concentration of excess carriers [26] when we impose a large difference between the potentials applied onto our gate electrodes, as necessary to engineer a deep lattice potential. As a result, excitons suffer additional collisions with free carriers, nevertheless, the photoluminescence is homogeneously broadened in this regime as well.

![FIG. 3: (a) Decay of the first-order time coherence for $V_0=3$ meV in log-log scale. Experimental data are displayed by the blue points while the red line shows an exponential decay with a characteristic time equal to 2.8 ps. The inset provides the same results in linear scale. (b) Standard deviation of the photoluminescence energy measured at the position of 3 adjacent lattice sites as a function of the lattice depth $V_0$. Each measurement is evaluated from an ensemble of 10 realisations all performed under the same experimental conditions. Experiments have all been performed at a bath temperature $T_b=340$ mK.](b)
tion we studied the variation of density fluctuations in the lattice sites as a function of the lattice depth. Indeed a Mott insulator is signalled by the same fixed number of particles in each site, so that density fluctuations are theoretically vanishingly small. For dipolar excitons density fluctuations are directly accessed by the energy of the photoluminescence which is governed by the strength of repulsive dipolar interactions between excitons [27-29]. Fig.3b compares the standard deviation of the photoluminescence energy \( \sigma_{E_X} \) for 3 neighbouring lattice sites as a function of \( V_0 \). Overall, it shows that the photoluminescence energy is highly stable in our experiments, since it varies by only around 80 \( \mu eV \). However, we do not observe any clear dependence of \( \sigma_{E_X} \) over \( V_0 \). In particular \( \sigma_{E_X} \) does not decrease while the lattice depth is ramped, and thus no sign of a Mott insulating phase is detected unambiguously.

In the deep lattice limit, the experiments displayed in Fig. 4 also allow us to study the time coherence of excitons trapped in the lattice sites. Indeed, in these measurements we ensured that one bright interference fringe coincided with one row of lattice sites. For each value of \( V_0 \) we then measured the intensity modulation along this particular bright fringe \( C_i \) and compared it to the bare modulation of the photoluminescence in real space \( C_m \) due to the localisation in the lattice sites. Thus, we directly extract the average amplitude of \( g^{(1)} \) for the photoluminescence radiated by the lattice sites only. It reads \( |g^{(1)}|=2(C_i-C_m)/((1-C_i)(1+C_m)) \), since the emission between the lattice sites does not yield any measurable interference signal for \( \tau=4.7 \text{ ps} \) (Fig.3.a). Figure 4 reveals then a revival of the interference contrast for \( V_0 \gtrsim 3 \mu eV \), the contrast increasing rapidly up to an amplitude similar to the one for vanishing \( V_0 \) for the deepest lattice depth. Accordingly, we deduce in the deep lattice limit that inside the sites dipolar excitons exhibit extended temporal coherence. This provides a signature for the build-up of microscopic quasi-condensates in the lattice sites [10, 19]. As shown in Fig.3, these coherent exciton droplets are independent one from another, i.e. with no defined phase relation. Let us then stress that we can not directly probe the spatial coherence of these droplets, since their spatial extension (\( \sim 1 \mu m \)) reduces to our spatial resolution. An extended time coherence then constitutes the only accessible fingerprint for the quasi-condensation of excitons trapped in the lattice sites.

To conclude, we have experimentally studied the quasi-condensate crossover for bilayer excitons confined in a lattice potential. We have observed that a quasi-condensate is formed when the barrier height is vanishing, but rapidly destroyed when it is increased to around 0.2 \( \mu eV \). We have shown that this behaviour is consistent with a renormalisation of the excitons effective mass by the lattice depth, which increases the critical phase-space density for the quasi-condensation crossover. On the other hand, in the

\[
V_0^* = V_0 + 2 \tilde{g}(n_{nc}^{(max)} - n_{nc}^{(min)}), \quad n_{nc}^{(max/min)}
\]

denoting the maximum/minimum for the density of non-condensed excitons leading to an increase in the effective lattice potential due to the strong mean field interaction. Their difference amounts to around 25% of the total population, as discussed previously. Moreover, Fig. 4e illustrates that for low barrier heights non-condensed excitons are mostly located at the position of the barrier, whereas their fraction is minimum at the position of the lattice sites in order to minimize the interaction energy. We therefore estimate that \( \tilde{g}(n_{nc}^{(max)}(r) - n_{nc}^{(min)}(r)) / \sim 0.25 \mu eV \) leading to \( m_X=1.5 m_X \) for \( V_0 \) around 0.2 \( \mu eV \). The effective critical phase space density for the excitons quasi-condensation then increases from 8 to 12 such that the actual exciton phase-space density in our studies is not any more high enough for forming a quasi-condensate, as experimentally observed.

**FIG. 4:** (a) First-order time coherence as a function of the barrier height \( V_0 \) in the weak lattice limit. The contrast measures the amplitude of the first-order time correlation function \( g^{(1)}(0,4.7ps) \). (b) Time coherence in the deep lattice limit. Here the contrast measures the amplitude of the intensity modulation along a bright interference fringe, i.e. along the weaker confining direction. The shaded areas display our instrumental resolution limited by the signal to noise ratio in these measurements. Experiments have all been performed at a bath temperature \( T_b=340 \text{ mK} \).
deep lattice limit, we have observed that an array of phase incoherent quasi-condensates localized at the lattice sites develops. However, in order to reach the Mott insulator regime, we estimate that the period of the lattice has to be decreased to less than around 1 µm, such that the on-site interaction becomes a few times larger than the thermal activation energy. Importantly, this regime is at close experimental reach, including in the limit of a few excitons per lattice sites.

Acknowledgments

Our work has been financially supported by the Labex Matisse, the Fondation NanoSciences (Grenoble), and by OBELIX from the French Agency for Research (ANR-15-CE30-0020). The work at Princeton University was funded by the Gordon and Betty Moore Foundation through the EPiQS initiative Grant GBMF4420, and by the National Science Foundation MRSEC Grant DMR 1420541.

[1] "Many-Body Physics with Ultracold Gases", Lecture Notes of the Les Houches Summer School (Eds. C. Salomon, G. V. Shlyapnikov, and L. F. Cugliandolo, 2010)
[2] M. Lewenstein, A. Sanpera, and V. Ahufinger, "Ultracold Atoms in Optical Lattices: Simulating quantum many-body systems" (Oxford University Press, 2012)
[3] M. Combescot, R. Combescot, F. Dubin, Rep. Prog. Phys. 80, 066401 (2017).
[4] P. Rivera, H. Yu, K. L. Seyler, N. P. Wilson, W. Yao, X. Xu, Nat. Nano. 13, 1004 (2018)
[5] C. Lagoin et al., submitted
[6] M. Remeika et al., Appl. Phys. Lett. 100, 061103 (2012)
[7] K. L. Seyler et al., Nature 567, 66 (2019)
[8] K. Tran et al., Nature 567, 71 (2019)
[9] C. Jin et al., Nature 567, 76 (2019)
[10] B. Laikhtman and R. Rapaport Phys. Rev. B 80, 195313 (2009).
[11] A. Filinov, N. V. Prokofev, and M. Bonitz, Phys. Rev. Lett. 105, 070401 (2010)
[12] A.A. High et al., Phys. Rev. Lett. 103, 087403 (2009)
[13] M. Beian et al., EuroPhys. Lett. 119, 37004 (2017)
[14] C. Schindler and R. Zimmermann, Phys. Rev. B 78, 045313 (2008).
[15] U. V. Butov, Superlattices Microstruct. 108, 2 (2017)
[16] Y. Y. Kuznetsova et al., Opt. Lett. 40, 589 (2015)
[17] C. Schindler and R. Zimmermann, Phys. Rev. B 78, 045313 (2008).
[18] B. Lailhtman and R. Rapaport Phys. Rev. B 80, 195313 (2009).
[19] K Berg-Sorensen, K Molmer, Phys. Rev. A 58, 1480 (1998)