Temporal coherence of spatially indirect excitons across Bose–Einstein condensation: the role of free carriers

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Abstract
We study the time coherence of the photoluminescence radiated by spatially indirect excitons confined in a 10 μm electrostatic trap. Above a critical temperature of about 1 K, we show that the photoluminescence is homogeneously broadened in the dilute regime, with a spectral width around 500 μeV that weakly varies with the exciton density. By contrast, the spectral width reduces by two-fold below the critical temperature and for experimental parameters at which excitons undergo a gray Bose–Einstein condensation. We find evidence showing that the photoluminescence temporal coherence is limited by interactions between excitons and a low concentration of residual excess charges, leading to a minimum photoluminescence spectral width of around 300 μeV in the condensed regime.

1. Introduction

In semiconductor physics, Bose–Einstein condensation of excitons, i.e. Coulomb bound electron–hole pairs, has received a large attention since original theoretical predictions were formulated in the 1960s [1–3]. Indeed, excitons are very attractive because their effective mass, much lighter than atoms, yields a cryogenically accessible critical temperature for the phase transition. However, the excitons electronic structure has hidden for a long time their condensation [4]. This property is well illustrated in widely studied GaAs heterostructures where excitons have four accessible ‘spin’ states, two optically dark states with spins (±2) and two optically bright ones with spins (±1). Bright excitons lying at a slightly higher energy than dark ones [5, 6], Bose–Einstein condensation leads to a macroscopic occupation of lowest energy dark states so that the condensate is not detectable easily. In fact, it is only in double quantum wells (DQWs) that a condensate can be imaged directly [7], when a small fraction (~20%) of bright excitons is possibly introduced coherently in a dark condensate by exciton–exciton interactions at sufficiently high densities (~10^{10}cm^{-2}) [7–9]. The condensate then becomes gray and is signaled by the weak photoluminescence it radiates.

Recently, we have reported a series of experiments combining the signatures expected for a gray condensate of indirect excitons [10] in GaAs DQWs. These signatures include first a (quantum) darkening of the photoluminescence below a critical temperature of about 1 K and for a narrow range of densities only (~2–3) × 10^{10} cm^{-2}) [11]. This darkening marks a quantum statistical occupation of dark excitons because the energy splitting between bright and dark states amounts to only a few μeV in DQWs [5, 6], i.e. at least 10 times less than the thermal energy at 1 K. We have then confirmed the gray nature of the condensation, i.e. the coherent coupling between dark and bright excitons. For that purpose we used spatial interferometry and showed that the weak photoluminescence radiated in the regime of quantum darkening exhibits macroscopic (quantum) spatial coherence [12], below the same critical temperature of about 1 K and for the same narrow density range.

Here, we scrutinize another fingerprint expected for exciton condensation, namely a steep increase of temporal coherence across the condensation threshold [9]. This behavior is understood by noting that in the
condensed phase optically bright excitons merely occupy \( k \sim 0 \) states, directly coupled to photons. At sub-Kelvin temperatures, the excitons coherence time is then limited by interactions between condensed and uncondensed \( (k \neq 0) \) excitons, as well as interactions between excitons and excess carriers which are strong in GaAs heterostructures [13]. Indeed, we note that exciton–phonon interactions are negligible in the sub-Kelvin regime [14]. Bose–Einstein condensation then has to increase the excitons coherence time because interactions between excitons of the condensate cannot contribute to collisional broadening due to energy and momentum conservation. In the following experiments, for a measured condensate fraction of about 80% [12] we show that the photoluminescence temporal coherence is increased by two-fold at threshold, the scattering rate of \( k \sim 0 \) bright excitons being then reduced by the same fraction. Our studies indicate that this amplitude is limited by the amount of free carriers, due to both the level of unintentional doping during the epitaxial growth of our heterostructure, and a possible exciton ionization due to residual in-plane electric fields reminiscent of electrostatic trapping for indirect excitons [15, 16].

2. Experimental procedure

We report time and spatially resolved interferometry to probe the temporal coherence of indirect excitons confined in a 10 \( \mu \)m trap. We study the same device as in [11, 12] in the regime where we have shown signatures of excitonic condensation. In this heterostructure, indirect excitons (IXs) are made by the Coulomb attraction between electrons and holes confined in two 8 nm GaAs quantum wells separated by a 4 nm AlGaAs barrier. The trap is engineered by using two metallic electrodes, a central 10 \( \mu \)m wide disk surrounded by an outer guard, both deposited at the surface of a field-effect device embedding the DQW. We apply the same potential on both electrodes \((-4.7 \text{ V})\), so that the 200 nm gap between them leads to a rectification of potential leading to an electrostatic barrier, i.e. a hollow-trap [12]. Moreover, note that IXs exhibit a long lifetime (~100 ns) in our studies and reach thermal equilibrium [17] in a mostly electrically neutral environment about 100 ns after optical injection, when a transient photocurrent \((\lesssim 100 \text{ pA D.C.})\) is evacuated. Indeed, we inject electronic carriers optically, using a 100 ns long laser excitation, repeated at a rate of 1.5 MHz and tuned on resonance with the direct exciton absorption of the two quantum wells (figure 1). IXs are thus obtained once electrons and holes have tunneled towards their minimum energy states, located in each quantum well. Let us then note that the laser beam covers most of the trap area.

To study the temporal coherence of bright IXs confined in the trap, we magnified the real image of the photoluminescence (PL) and sent it to a Mach–Zehnder interferometer. The PL is divided equally between a fixed and a mobile arm, and the two beams are superposed at the output of the interferometer, i.e. without any spatial displacement. Horizontal interference fringes are induced by a vertical tilt angle between the two outputs, of about 25°. To evaluate the temporal coherence of the PL, we vary the longitudinal path length difference, i.e. we translate one arm by \( \delta L \). The two outputs are then shifted temporally by \( \delta t = 2\delta L/c \), where \( c \) is the speed of light in air, the path length difference being actively stabilized with \( \sim 20 \text{ nm} \) precision. At the output, the contrast of the interference pattern measures directly the coherence time \( \tau_c \) of optically bright excitons. Indeed, the interference visibility is given by the modulus of the first-order correlation function \( |g^{(1)}(r, \delta t)| \propto \left| \langle \psi_x^*(r, t) \psi_y(r, t+\delta t) \rangle \right| \). Here, \( \psi_x(r, t) \) is the photoluminescence field which reflects the wave function of bright excitons with an in-plane momentum \( k \approx 0 \) since the exciton–photon coupling is linear close to resonance [18]. Moreover, \( \langle \cdot \rangle \) denotes the time averaging and \( \mathbf{r} = (x, y) \) reads the coordinate in the plane of the electrostatic trap. For a homogeneously broadened gas, the \( g^{(1)} \)-function is exponentially decaying with \( \delta t \), at a rate inversely proportional to the \( k = 0 \) excitons coherence time \( \tau_c \). In a thermal regime above a few Kelvin, \( \tau_c \) is controlled by exciton–phonon interactions [14]; however, below a few Kelvin, we expect a dramatic change because these interactions become inefficient.

3. Spectral fluctuations and excitons coherence time

Unlike studies performed with atomic gases [19], our experiments require an accumulation of several millions of single-shot realizations, during around 10–20 s, each realization starting by a loading laser pulse after which the PL is interferometrically analyzed (figure 1). This approach is inevitable due to the weak photoluminescence signal radiated from the trap, but it constitutes a severe limitation. Indeed, we observe large spectral fluctuations, with a characteristic timescale of about 10 s, as shown in figure 1(a) at a bath temperature \( T_b = 330 \text{ mK} \) and for a trapped exciton density \( n_X \) about 2.5 \( \times 10^{10} \text{ cm}^{-2} \). Particularly, figure 1(a) displays two spectra measured successively; a first spectrum consists mostly of a central narrow line, with a width limited by the spectrometer resolution. By contrast, another spectrum obtained for the same experimental settings can be about 1 meV broad with no recognizable shape (dashed line in figure 1(a)).
Figure 1. Our experimental sequence starts with a 100 ns laser excitation loading indirect excitons in the trap. The reemitted photoluminescence is then analyzed in a 5 ns long time interval, at a given delay to the end of the loading laser pulse so that the density in the trap is thus varied. The sequence is repeated at 1.5 MHz.

(a) Spectrally resolved PL emitted across the center of the trap at $T_b = 330$ mK. Measurements are realized during a 20 s long acquisition and for an exciton density $n_X \sim 2.5 \times 10^{10}$ cm$^{-2}$. The sharp line is fitted with a Lorentzian shown by the dotted line (FWHM $= 430 \mu$eV), whereas the dashed line is anomalously broad ($\sim 1$ meV).

(b) Distribution of interference contrasts for a sample of 100 acquisitions measured in the same experimental conditions as (a). The best 10% of these realizations is highlighted in yellow, a minimum contrast ($\sim 15\%$) being set by the signal to noise ratio at the detection. High contrast (c) and lower contrast (d) interference patterns measured for the same experimental conditions as in (a), but during a 12 s long acquisition, the interferometer path length difference being set to $\delta t = 1.8$ ps. (e)–(f) Interference signal taken along the center of the trap ($x = 5$ $\mu$m) for the measurements shown in (c) and (d), respectively. The interference contrast is evaluated by computing the modulation amplitude at the center of the trap, i.e. for $x = 5$ $\mu$m and $y$ between 3 and 7 $\mu$m, considering that the background amounts to 10 counts in our experiments.
We now study the variation of the photoluminescence coherence time as a function of the exciton density. On the one hand, this property is useful since it allows us to extract the density of excitons in the trap $n_X$, indirect excitons experiencing repulsive dipolar interactions that yield a blue shift of the photoluminescence energy scaling as $\Delta E_{\text{dip}}$ [20, 21], where $\Delta E_{\text{dip}} \sim 0.5 \text{ meV/m}^{10}\text{cm}^{-2}$. On the other hand, indirect excitons interact strongly with free carriers, due to either the laser induced photocurrent or to the level of residual doping during the growth of our heterostructure. In both cases, the exciton-free carrier interactions provide a direct source for spectral diffusion because the local concentration of excess carriers cannot be controlled and a priori varies randomly during our measurements. To minimize this effect, we typically analyze the PL long after extinction of the laser excitation, i.e. when the photocurrent is damped. Nevertheless, even a density of excess carriers small compared to $n_X$ can yield a significant electrostatic noise, because in GaAs quantum wells the interaction strength between excitons and free carriers is about 10 times larger than the one between excitons [13].

Spectral fluctuations are also evidenced interferometrically, with actually better dynamics since the acquisition time is reduced from about 20 to 10 s once the PL is not filtered spectrally. For fixed experimental conditions at $T_b = 330 \text{ mK}$, figure 1(c) shows that either a clear interference pattern can be observed across the center of the trap, or a rather blurred pattern (figure 1(d)). For these experiments where the path length difference is set such that $\delta t = 1.8 \text{ ps}$, in the former case the contrast along the center of the trap ($x = 5 \mu\text{m}$) amounts to 35% (figure 1(e)) whereas in the latter case it reduces to about 23% (figure 1(f)) manifesting an increased spectral width. Such variation is quantified in figure 1(b) where we display the interference visibility measured at $T_b = 330 \text{ mK}$ for a statistical ensemble of 100 experiments performed all under the same conditions. We note that the contrast exhibits an asymmetric dispersion, from 17% to 37%, marked by a tail at high contrasts (30%–37%) and a sharp growth at low contrasts (17% to 22%). While we cannot quantify the origin of such asymmetry, it indicates that the spectral diffusion is complex and not simply due to a Gaussian noise process. Based on the calibration of our interferometer, we nevertheless deduce that the photoluminescence exhibits a spectral width in average lower than 700 $\mu\text{eV}$ and that for the most regular electrostatic environment, e.g. for a lowest concentration of excess free carriers, the spectral width can become as narrow as 300 $\mu\text{eV}$.

4. Excitons temporal coherence

We now study the variation of the photoluminescence coherence time as a function of the exciton density $n_X$ and bath temperature $T_b$. To limit the influence of spectral fluctuations over our conclusions, we restrict our analysis to the realizations with highest visibility for each experimental settings. Precisely, we average the contrasts measured for the best 10% of experimental realizations, post-selected within a statistical ensemble of a minimum of 50 measurements for each experimental conditions (see for instance the orange colored region in figure 1(b)). Thus, we aim at minimizing the amount of measurements where inhomogeneous broadening plays a significant role. Instead, we focus onto experiments where indirect excitons are confined in the most regular trapping potential, combined to a minimum amount of excess charges interacting with them. Post-selecting these optimum conditions is actually necessary to study the excitons gray condensation, since we have shown in [12] that a gray condensate is only obtained for such optimized experimental settings and cannot form otherwise.

Let us first discuss the dilute regime, i.e. an exciton concentration $n_X$ of about $2.5 \times 10^{10} \text{ cm}^{-2}$ obtained for a 170 ns delay to the extinction of the loading laser pulse. Figure 2 quantifies the variation of the interference contrast measured along the vertical axis at the center of the trap ($x = 5 \mu\text{m}$), as a function of the time delay $\delta t$ introduced between the two arms of the interferometer. We note that the interference visibility follows a mono-exponential decay $e^{-|\delta t|/\tau}$, at every bath temperature $T_b$ below 2.36 K. According to the Wiener–Khintchine theorem [22], the photoluminescence spectrum then consists of a single Lorentzian. This shows that we study a homogeneously broadened gas which constitutes a major improvement compared to previous works [23–28]. To the best of our knowledge, such a single Lorentzian PL profile had never been observed for a gas of indirect excitons. For the measurements shown in figure 2, the photoluminescence spectral width $\Gamma = 2\hbar/\tau$, reaches then 550 $\mu\text{eV}$ at $T_b = 2.4 \text{ K}$. This value is nevertheless $\sim 300 \mu\text{eV}$ greater than the theoretically expected linewidth due to exciton–exciton interactions [20], indicating that these latter are not the only mechanism governing the homogeneous broadening.

We now discuss the photoluminescence temporal coherence as a function of the exciton density. Experimentally, this is achieved simply by varying the delay to the loading laser pulse. Indeed, the exciton density decreases monotonously in the trap due to radiative recombination, and let us stress that in our studies the total density of excitons $n_X$, including both bright and dark states, is the same at a given delay for every bath

$$n_X = \frac{\hbar c}{\Delta E_{\text{dip}}} = 500 \text{ Debye}.$$
temperature below 3 K \(1\). For a fixed delay \(\delta t = 1.8\) ps, figure 3 shows that three regimes emerge: (i) in the most dilute regime, i.e. for longest delays to the end of the loading laser pulse, the interference contrast does not vary with \(T_b\). This behavior indicates that the gas is localized, as expected since in this range of densities the blue shift of the photoluminescence, i.e. the excitons mean interaction energy, does not exceed the electrostatic disorder (\(\sim 500\) \(\mu\)eV). At higher densities (regime (ii)) the excitons mean-field energy compensates the strength of potential disorder so that dipolar repulsions delocalize the exciton gas. Then, the interference contrast varies strongly with \(T_b\): while the visibility keeps a value similar to regime (i) for \(T_b \approx 1.2\) K, it increases steeply at \(T_b \approx 330\) mK, from about 24\% to 30\% when \(n_X \gtrsim 2 \times 10^{10}\) cm\(^{-2}\). Increasing further the density, \(n_X \gtrsim 3.5 \times 10^{10}\) cm\(^{-2}\) in regime (iii), leads to a breakdown of this behavior and the interference contrast depends again weakly on the bath temperature.

From the visibility measured for \(\delta t = 1.8\) ps, we can deduce the spectral width \(\Gamma = 2\hbar/\tau_c\) in regime (ii) where we have shown in figure 2 that the photoluminescence is homogeneously broadened. At \(T_b = 1.2\) K, figure 3 shows that the homogeneous spectral width is about 500 \(\mu\)eV at 1.2 K and varies weakly with the exciton density \(n_X\). At \(T_b = 2.4\) K, \(\Gamma\) is slightly larger (\(\sim\)600 \(\mu\)eV), possibly due to the increased role of acoustic phonons, while it also varies weakly with \(n_X\). Such weak variations with \(n_X\) are theoretically expected \(20, 29\), since interactions between indirect excitons vary \(\Gamma\) by less than 100 \(\mu\)eV in the density range explored here, i.e. by less than our instrumental precision. However, the magnitude of \(\Gamma\) lies about 300 \(\mu\)eV above its theoretical expectation. We attribute this discrepancy as a manifestation of the interaction between indirect excitons and

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure2.png}
\caption{Interference visibility measured along the center of the trap as a function of the time delay introduced in the interferometer. Measurements taken between \(T_b = 330\) mK (red circles) and 2.36 K (blue squares) all follow single exponential decays (solid lines). The insert displays the data at \(T_b = 330\) mK with an extended abscissa, fitted by a single exponential with a time decay of 4 ps.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure3.png}
\caption{(a) Interference contrast \(|g^{(1)}|\) measured along the center of the trap as a function of the exciton density, \(n_X\). Measurements at \(T_b = 330\) mK, 1.2 and 2.4 K, are displayed in red, green and blue, respectively. Regions (i)–(iii) are discussed in the main text.}
\end{figure}
residual excess charges, although we minimize the concentration of these latter in regime (ii) since the delay to the laser excitation exceeds 120 ns. Let us then note that interactions between excitons and free carriers yields a Lorentzian broadening of the photoluminescence [30].

5. Excess free carriers and collisional broadening

To study the role of excess carriers in the DQW, let us first estimate the fraction of free carriers induced by the laser excitation. For that purpose, we consider the dynamics of the photocurrent, which average $I$ amounts to 100 pA, as measured in a 20 ms time interval with the source-measure device used to polarize the exciton trap. Moreover, we have previously shown that the photocurrent has a transient time of the order of 70 ns [11]. In this case we deduce that $I = I_m T_p$, where $I_m$ is the maximum current generated at the termination of the laser excitation, while $f$ and $T_p$ denote the 1.5 MHz repetition rate of our sequence and the length of the loading pulse, respectively. This expression for $I$ directly converts into $n_m$, the maximum density of excess carriers induced at the termination of the laser excitation, since $n_m = I_m/(e f r^2)$, where $e$ is the electron charge. For the radius $r = 5 \mu m$ of our trap electrode, we conclude that $n_m = 3.5 \times 10^{10} \text{ cm}^{-2}$ so that the density of photo-induced excess carriers reduces exponentially to a few $10^8 \text{ cm}^{-2}$ in regime (ii) of figure 3. On the other hand, excess carriers also result from the level of unintentional doping during the growth of our heterostructure. Indeed, the concentration of impurities for the molecular beam epitaxy of our sample is at most $10^{15} \text{ cm}^{-3}$, yielding a density of free carriers of about $10^8 \text{ cm}^{-2}$ for two 8 nm wide quantum wells. Unintentional doping contributes then much more significantly to the concentration of free carriers interacting with excitons long after the loading laser pulse. Finally, let us note that free carriers can also result from exciton ionization at the boundary of our electrostatic trap. Indeed, we use the rectification of the electrostatic potential at the interface between the two surface electrodes to define the trap area. This rectification leads to an electric field component normal to the quantum wells, thus inducing a potential barrier for indirect excitons. However, it also leads to field components in the plane of the quantum wells [15]. Although these latter are normally small for our heterostructure design, we cannot exclude that some fraction of excitons is ionized in the outer rim of the trap, in which case excess holes would remain confined under the trap area [31].

To estimate the impact of excess carriers onto the photoluminescence spectral width, we compare the strength of collisional broadening due to exciton–exciton interactions, in single and DQWs. In the former case, experiments by Honold et al [13] have shown that the photoluminescence spectral width is increased by about 100 $\mu$eV when the exciton density grows from $10^{10}$ to $2 \times 10^{10} \text{ cm}^{-2}$. This collisional strength is of the same order as the calculations made by Zimmermann for indirect excitons in DQWs [20], neither of which contradict our observations where collisional broadening between indirect excitons is not observed with our $\sim 100 \mu$eV resolution (regime (ii) in figure 3). This reasonable agreement suggests that in both geometries exciton–exciton interactions have a comparable strength. Let us then extend to DQWs the strength of interactions between excitons and excess carriers measured in single quantum wells, which is 8 times larger than interactions between excitons [13]. To account for the $300 \mu$eV mismatch between the homogeneous linewidth predicted in [20] and our measurements at $T_p \gtrsim 1.2 \text{ K}$ in regime (ii), we deduce that a free carrier density of about a few $10^9 \text{ cm}^{-2}$ is sufficient. Remarkably, this concentration is of the order of the level of unintentional doping thus suggesting that this limitation is potentially responsible for about $300 \mu$eV of the measured photoluminescence spectral width. Finally, let us note that by considering $n_m$ and the level of unintentional doping, together with the strength of interactions between excitons and excess carriers measured by Honold et al [13], we deduce that the photoluminescence spectral width at the termination of the laser excitation is mostly governed by interactions between excitons and photo-injected free carriers, and amounts to about 1 meV. This value is in good agreement with our observations [11], thus supporting further the previous assumptions. In fact, we attribute the rapid and temperature independent increase of $\Gamma$ in regime (iii) to excitons-free carriers interactions. Indeed, the latter regime corresponds to delays ranging from 120 to 20 ns after extinction of the loading laser pulse. While the exciton density is increased by analyzing shorter delays, the concentration of excess carriers is also increased since the transient photocurrent is not fully suppressed and therefore $\Gamma$ increases in a temperature independent fashion governed by excitons-free carriers scatterings.

6. Threshold increase of excitonic coherence

The most striking feature of figure 3 is certainly the threshold increase of the interference contrast at $T_p = 330 \text{ mK}$ and $n_X \sim 2 \times 10^{10} \text{ cm}^{-2}$, so that $\Gamma \sim 300 \mu$eV and remains rather constant until $n_X \sim 3.5 \times 10^{10} \text{ cm}^{-2}$. Remarkably, this density range coincides with the one where we have demonstrated a gray Bose–Einstein condensate for identical experimental conditions [12]. It is then natural to attribute the observed threshold as the result of the exciton condensation. Furthermore, $\Gamma \sim 300 \mu$eV matches the...
somehow not very surprising since longer coherence times are expected from exciton concentration long after a laser excitation, about $10^9$ cm$^{-2}$, limited to around 4 ps, due to interactions between indirect excitons and excess carriers. These latter have a low photoluminescence temporal coherence. Our experiments reveal that the condensate has a coherence time $\tau_c$.

We have shown that the gray condensation of indirect excitons is associated to a threshold increase of the exciton density $n_X \sim 2.5 \times 10^{10}$ cm$^{-2}$.

As shown in figure 4(a), a rapid increase is revealed below a bath temperature of about 800 mK, $T_b$, passing from about 2 ps at 1 K to 4 ps at 330 mK. To measure more precisely the increase of $\tau_c$, we performed an additional experiment where the interference contrast was evaluated for $\delta t = 1.8$ ps only. Thus, the stability of our measurements was largely increased since the experiments could be performed during a single day, compared to 3 d for the data displayed in figure 4(a). As shown in figure 4(b) we then recovered that two regimes emerge: for $T_b \geq 1.2$ K, $|g^{(1)}(\mathbf{r}, \delta t)|$ varies weakly, it is about 22% corresponding to $\tau_c \sim 2$ ps. By contrast, below $T_c \sim 1.2$ K, the interference contrast increases abruptly to 32% at $T_b = 330$ mK, manifesting that $\tau_c$ is doubled at sub-Kelvin bath temperatures. Let us then stress that this value of $T_c$, reproduces quantitatively the critical temperature measured for the appearance of quantum spatial coherence for the same exciton density in the trap [12]. This strongly supports our interpretation that the increase of excitonic temporal coherence is due to their quantum condensation in the trap.

7. Conclusion

We have shown that the gray condensation of indirect excitons is associated to a threshold increase of the photoluminescence temporal coherence. Our experiments reveal that the condensate has a coherence time limited to around 4 ps, due to interactions between indirect excitons and excess carriers. These latter have a low concentration long after a laser excitation, about $10^9$ cm$^{-2}$, i.e. of the order of the level of unintentional doping during the epitaxial growth of our heterostructure. In this situation we do not find any indication for the formation of trions, i.e. bound states between excitons and excess carriers, neither in the normal nor in the condensed regime. Furthermore, from our observations we anticipate that reducing the density of free carriers to a minimum value of a few $10^9$ cm$^{-2}$ would lead to a maximum time coherence of several tens of ps for gray condensates. This may be accessed following the technology developed for electrically connected bilayer heterostructures [32], i.e. by using additional gate electrodes contacting directly each quantum well. Moreover, our experiments indicate that quantum condensation is excluded for densities of free carriers exceeding $\sim 5 \times 10^9$ cm$^{-2}$, as expected theoretically [33]. In every case, we note that the coherence time of a gray condensate remains one to two orders of magnitude smaller than for quantum fluids of polaritons [34]. We

![Figure 4](image-url)

Figure 4. (a) Coherence time of the photoluminescence $\tau_c$ as a function of $T_b$, deduced by fitting the exponential decay of $|g^{(1)}(\mathbf{r}, \delta t)|$ at every bath temperature. (b) Interference visibility measured at the center of the trap for $\delta t = 1.8$ ps. The inset displays the value of $\tau_c$ deduced from these measurements. All experiments have been carried out for an exciton density $n_X \sim 2.5 \times 10^{10}$ cm$^{-2}$.
attribute this difference as the manifestation of the matter-like nature of exciton condensates which are then very sensitive to their electrostatic environment.

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