Deterministic assembly of a charged quantum dot - pillar cavity device.

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Developing future quantum communication may rely on the ability to engineer cavity-mediated interactions between photons and solid-state artificial atoms, in a deterministic way. Here, we report a set of technological and experimental developments for the deterministic coupling between the optical mode of a micropillar cavity and a quantum dot trion transition. We first identify a charged transition through in-plane magnetic field spectroscopy, and then tune the optical cavity mode to its energy via in-situ lithography. In addition, we design an asymmetric tunneling barrier to allow the optical trapping of the charge, assisted by a quasi-resonant pumping scheme, in order to control its occupation probability. We evaluate the generation of a positively-charged quantum dot through second order auto-correlation measurements of its resonance fluorescence, and the quality of light-matter interaction for these spin-photon interfaces is assessed by measuring the performance of the device as a single-photon source.

The efficient interfacing of single-photons to natural or artificial atoms drives advances in solid-state photonic quantum networks.1 Strong atom-photon interactions cannot be achieved in homogeneous electromagnetic environments, hence requires to be enhanced. This can be obtained by inserting the atom in an optical resonator, thus resulting in a cavity quantum electrodynamics (cQED) device2,3.

An important application for such devices is the emission of single photons: a good single-photon source should deterministically produce pure and indistinguishable single photons4,5. Among other systems, quantum dots (QDs) are excellent candidates, able to emit highly-pure single-photons6,7 with near-unity indistinguishability8–12. Moreover, their collection efficiency can be greatly increased by inserting them into high quality-factor cavities8,11–16. Both neutral and charged quantum dot transitions can be used as efficient single-photon sources, yet higher brightness is commonly achieved with a charged quantum dot15.

A cQED device can also act as an efficient receiver of incoming photons3, where the state of an artificial atom controls, or is controlled by, a single photon. Such interfacing allows using the artificial atom as a quantum memory to store the photonic state, and to realize various atom-photon, atom-atom or photon-photon gates17,18. This can be performed by addressing the exciton qubit in a neutral QD, and use the giant optical non-linearity of the cQED device to develop a single-photon router19. Another possibility is polarization encoding, using the interaction with a neutral QD to induce a large polarization rotation on the reflected photons20–22. However, a neutral QD has only one single ground state and therefore, does not have degrees of freedom to be manipulated23.

To use cQED devices as efficient light-matter interfaces, one can thus use the spin degree of freedom of a singly-charged QD, that is interfaced with the cavity mode23–25. Promising experimental results were already achieved either using a spin to either control the path of a single photon26 or to induce cavity-enhanced phase-shifts or giant polarization rotations27–31.

However, so far the deterministic fabrication of such singly-charged cQED devices with nano-engineered cavities is still to demonstrate. Such a deterministic realization indeed needs to fulfill challenging requirements at the same time. The first one is the deterministic coupling of a singly-charged QD state to a cavity mode. A second challenge is to deliberately prepare the QD in the desired charge state, consisting in either an electron or an hole in excess. Finally, the cavity must allow the photons to be efficiently injected and collected after having interacted with the QD. In this respect, pillar microcavities allow for both high efficient injection32–34 and collection of photons15.

Here, we demonstrate the deterministic realization of efficient singly-charged cQED devices. By deterministically interfacing a single positively-charged QD to a micropillar cavity, we can both inject and collect single-photons efficiently. This is achieved by pre-selecting the trion transition through in-plane magnetic field photoluminescence, and then spectrally coupling it to the cavity mode via the in-situ lithography technique35,36. A single hole is efficiently trapped inside the QD thanks to an engineered asymmetric band structure, that hinders the tunneling of the hole out of the QD, and thanks to an optical injection technique. The efficiency of this injection technique is then evaluated via intensity correlations of the QD resonance fluorescence. Finally, the properties of the device are evaluated, operating as a single photon source.

This paper is organized as follows. In Sec. I, we present the assymetric band structure design. The deterministic coupling of the cavity to the trion transition is discussed in Sec. II. Then, we present the optical trapping scheme in Sec. III, and evaluate its performances in Sec. IV. Fi-
nally, the properties of the devices operating as a single-photon source are discussed in Sec. V.

I. ASYMMETRIC BAND STRUCTURE ENGINEERING.

The fabricated samples consist in a λ-cavity, surrounded by two distributed Bragg reflectors (GaAs and Al$_{0.9}$Ga$_{0.1}$As, with 14 and 28 pairs for the top and bottom mirrors respectively), thus forming a pillar microcavity. The reduced number of layer pairs in the top mirror leads to intracavity photons escaping the cavity from the top with a probability $\eta_{\text{top}} = 85 \pm 5\%$. The λ-cavity consists in a GaAs layer which embeds a single InGaAs QD and a tunneling barrier layer of Al$_{0.1}$Ga$_{0.9}$As, whose role is detailed afterwards.

Figure 1. (a) Simulated electromagnetic field intensity inside the micropillar structure. The electromagnetic field is confined at the cavity layer position (green: GaAs layers, yellow: Al$_{0.9}$Ga$_{0.1}$As layers, purple: Al$_{0.1}$Ga$_{0.9}$As tunneling barrier). (b) Scanning electron microscopy image of a sample. (c) QD + doping structure: the n-doping and p-doping region tilt the forbidden band so that fluctuating charges remain far from the QD, therefore stabilizing the electrical fluctuations. The tunneling barrier reduces the hole tunneling rate.

Fig. 1(a) displays the simulated electromagnetic field intensity plotted as a function of the position in the structure (due to its small width, the QD layer was neglected in these simulations). The micropillar structure is designed to confine the electromagnetic field with a maximum intensity at the QD layer position in order to maximally interact with a single QD.

In addition, the micropillar structure contains a p-i-n junction with a doping density that gradually decreases while approaching the cavity (negative doping for the bottom mirror and positive doping for the top mirror). It is connected to an electrically-contacted diode through four ridges and a circular frame as shown in Fig. 1(b). A well known technique to control the charge state of QDs is via the bias voltage, which brings the Fermi energy close to a charged QD state. For QDs coupled to microcavities, however, the free-carriers of the doped regions can absorb the intracavity photons leading to additional cavity losses. In this respect, the doped regions should be positioned far from the QD position, where the electromagnetic field of the confined cavity mode is not intense. This hinders the obtention of a Fermi level very close to the singly-charged QD state. Another strategy to optically confine a single hole into the QD is proposed here, based on the original ideas proposed in Ref. 41.

The energy levels of the QD and its near environment are schematized in Fig. 1(c), which sketches the energy bands as a function of the vertical position in the cavity layer. The doped regions tilt the forbidden bands such that an electron confined inside the QD can tunnel out quickly ($< 1\mu$s). However, thanks to the 20nm-thick Al$_{0.1}$Ga$_{0.9}$As tunneling barrier that is positioned 10nm above the QD layer, the hole tunnels out with a much longer timescale than the electron, thus facilitating the confinement of a single hole: if an electron-hole pair (also called an exciton X) is optically generated, the electron escapes the QD at a much faster rate than the hole. Such an asymmetrical design should thus at the same time increase the hole confinement time and favor the single hole QD state over the electron QD state. A strategy of optical confinement of the hole inside the QD is proposed here, based on the original ideas presented in Ref. 41.

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II. DETERMINISTIC TRION-CAVITY COUPLING.

Because the doping is situated 200nm away from the QD layer, the identification technique of the QD states based on the photoluminescence under applied voltage is not working. The coupling between the trion transition and the optical mode of the cavity is achieved by analyzing the photoluminescence of the trion transition under a transverse magnetic field. The identified trion transition can then be coupled both spatially and spectrally to the QD via in-situ lithography.

The trion transitions with and without magnetic field are sketched respectively in the top and bottom panels of Fig. 2(a). At zero external magnetic field (bottom panel of Fig. 2(a)), a trion radiatively decays into the hole state by emitting a circularly-polarized photon with either a right-handed or a left-handed helicity. There
are therefore two trion transitions that are energy degenerate. When an in-plane magnetic field is applied, it induces a Zeeman splitting between the hole states and the trion states, and it also modifies the system eigenstates and thus the optical selection rules, as illustrated in the top panel of Fig. 2(a). For a high in-plane magnetic field ($B_x > 1T$ typically), a trion decays into a hole by emitting a single photon with 4 combinations of linear polarizations and energies. This feature is a signature of a trion state and is used experimentally to identify it on a planar cavity sample in Fig. 2(b).

The bottom panel of Fig. 2(b) shows a typical photoluminescence spectrum obtained on a QD in a planar cavity sample, under non resonant excitation ($\lambda_{NR} = 850nm$) and without any applied magnetic field. There are three dominant transitions observed, at 925.1nm, 925.3nm and 925.7nm: the lower intensity transition (at 925.5nm) is related to another QD next to the one under study. The spectrometer resolution (25pm) does not permit to resolve the exciton fine structure\textsuperscript{42}. Consequently, from such spectrum, it is impossible to identify the QD states.

When an in-plane magnetic field is applied, the QD transitions start to split. This can be observed in the middle panel of Fig 2(b), which displays the evolution of the photoluminescence spectrum under an increasing in-plane magnetic field intensity. At $B_x = 4T$ and above, all the observed peaks are split in two transitions. At high magnetic field, all the observed QD transitions are also blue-shifted, due to the diamagnetic shift\textsuperscript{13}. The trion transition splits into four different transitions; however, the limited spectrometer resolution impedes the visualization of such effect.

To overcome this difficulty, the photoluminescence spectrum is analyzed in polarization. The top panel of Fig. 2(b) displays the polarization-resolved photoluminescence for the same QD at $B_x = 4T$. The red (green) curve corresponds to the horizontally-polarized (vertically-polarized) photoluminescence. The horizontal (vertical) polarization peak centers are highlighted by red (green) vertical lines.

Let us first consider the higher wavelength transition (925.7nm at $B_x = 0T$) from the bottom panel of Fig. 2(b). The magnetic field scan and the polarization analysis show that it is split in two transitions with orthogonal linear polarizations. This is the expected behavior for an exciton\textsuperscript{42}.

As can be seen in the middle and top panels of Fig. 2(b), the lower wavelength transition (925.1nm at $B_x = 0T$) analysis shows that it is split in four transitions with different energies and polarizations: the highest and lowest have the same horizontal (H) polarizations and are orthogonal to the two center vertically-polarized (V) transitions. This state is therefore identified as a trion, as this is the behavior expected from the polarization selection rules of Fig. 2(a).

The center transition (925.3nm at $B_x = 0T$) analysis is also split in four transitions with similar polarizations as the QD trion. However, there is a clear asymmetry in the photoluminescence intensity for these four transitions. This feature may be explained by another excitonic...
state that for example can be a bi-trion state $X^{2+}$ (two holes and an electron-hole pair)\textsuperscript{44,45}.

Thus, the analysis in polarization of the QD photoluminescence under high transverse magnetic field allows to identify the trion transition wavelength. This transition can then be coupled to the fundamental cavity mode of a pillar microcavity by the in-situ lithography technique\textsuperscript{35,36}.

After the realization of electrically-contacted QD-micropillar devices, the trion-cavity spectral matching is evaluated by repeating the procedure of trion identification under an in-plane magnetic field. The trion transition at $B_x = 0$T is then compared to the cavity mode energy to verify that they are indeed coupled: the result is displayed in Fig. 2(c). The bottom panel represents the photoluminescence at $B_x = 0$T, displaying discrete peaks: the most intense peak corresponds to the QD transition that is coupled to the fundamental cavity mode. The middle panel of Fig. 2(c) shows the photoluminescence of a cavity-coupled QD under magnetic field. The top panel of Fig. 2(c) represents the photoluminescence collected for an external transverse magnetic field of approximately 6$\text{T}$ in H and in V polarizations. The QD transition that is in resonance with the fundamental cavity mode is characterized by four Zeeman-split transitions and is therefore a trion.

We showed that identifying the trion transition before proceeding to an in-situ lithography allows to deterministically couple the hole charge state to intracavity photons. Therefore, when a hole is trapped in the QD, its spin degree of freedom will be interfaced with photons. The next step is to make sure that a single hole is indeed confined in the QD with high probability.

III. OPTICAL CONTROL OF THE SINGLE-HOLE STATE IN THE QUANTUM DOT.

In this section, we propose an efficient optical injection scheme to load the extra charge in the trion state. In general, the stable QD state is the crystal ground state, here denoted as the neutral QD whose energy levels are represented in the left part of Fig. 3 (a). An electron-hole pair can be generated by resonantly exciting the exciton ($\omega_X$) or other discrete transitions (such as "p-shell" transitions or phonon-assisted transitions) here denoted as quasi-resonant (QR) transitions with energy $\omega_{QR} > \omega_X$.

In this work, an electron-hole pair is generated by a quasi-resonant CW laser at a wavelength $\lambda_{QR} = 901$nm as depicted in the left panel of Fig. 3(b). The exact QD state that is excited can correspond either to a p-shell excitation or to the generation of an s-shell exciton assisted by a longitudinal optical phonon\textsuperscript{46,47}.

The electron and the hole non-radiatively decay (in typically less than 100ps\textsuperscript{48}) into the QD and eventually form an exciton as shown in the middle panel of Fig. 3(b). If the exciton radiatively recombines by emitting a photon, the QD is returning to its ground state. This generation-recombination cycle keeps going until an excited electron eventually tunnels out towards the n-doped electrical contact before the radiative recombination occurs. Due to the Al$_{0.1}$Ga$_{0.9}$As barrier, the remaining hole is confined in the QD for a longer time, during which the QD is in the desired positively-charged state. The confined charge induces a strong modification of the electric environment by Coulomb interaction, modifying the energies of the discrete levels of the QD, as illustrated in the right hand side of Fig. 3(a). Therefore the trion transition does not have the same energy as the exciton transition ($\omega_X \neq \omega_X^{+}$). Similarly, the excess charge modifies all the other discrete energy levels and thus all the quasi-resonant transitions. Contrary to an above-band excitation scheme, which constantly generates electron-hole pairs, $\omega_{QR}$ does not correspond to any optical excitation when a hole is confined. This limits the

![Figure 3. (a) Sketch comparing the energy levels of a neutral QD and a singly-charged QD. The resonant trion transition energy $\omega_X^{+}$ is different from the exciton transition energy $\omega_X$ ($\omega_X^{+} \neq \omega_X$) and the p-shell transition energies of a charged QD are also different from a neutral QD: $\omega_{QR} \neq \omega_{QR}$. (b) Hole trapping scheme using quasi-resonant excitation. Left panel: the QD stable state is the crystal ground state. A quasi-resonant laser creates an electron-hole pair by exciting a p-shell transition ($\omega_{QR}$). Middle panel: If the electron tunnels out the QD before radiative recombination with the hole, it generates of a single hole QD state. Right panel: the QD can then be resonantly excited by a laser with energy $\omega_X^{+}$ corresponding to the trion transition.](image-url)
risk of confining more than one charge in the quantum dot.

The single hole thereby generated can then be optically manipulated by a second laser in resonance with the trion transition, i.e. at the energy $\omega_{X^+} (\neq \omega_X)$ with typical power $P_{X^+} \approx 0.1 - 3 \text{nW}$, as schematized in the right panel of Fig. 3(b). Because the exciton and the trion transition have different energies, this second laser does not intervene in the hole trapping procedure.

Instead of quasi-resonant excitation, another possibility to trap a hole would have been to use a laser in resonance with the QD exciton, $\omega_{\text{laser}} = \omega_X$. However, in the context of cQED with the cavity centered on the trion transition, the exciton is not in resonance with the cavity. Therefore, resonant exciton excitation requires a high laser intensity to inject photons into the cavity at a wavelength still close enough to the trion transition. Such an intensity is challenging to suppress on the reflected path, even when using a cross-polarized setup to collect only the single photons emitted by the resonance fluorescence of the QD.

**IV. PERFORMANCE OF THE HOLE TRAPPING SCHEME**

The hole trapping scheme performances, namely the single-hole occupation probability and the single-hole trapping time, are now evaluated by observing the resonance fluorescence signal on multiple devices. To suppress the resonant laser and only collect the emitted single photons, we use a cross-polarized setup where the quantum dot is resonantly excited with a horizontally-polarized resonant laser. The resonance fluorescence is collected through a vertical polarizer that filters out the resonant laser (the quasi-resonant laser is spectrally filtered). The 15-ps pulse laser in resonance with the trion transition induces Rabi oscillations between the hole and the trion state$^{10,49,50}$. In the following, we refer to the laser pulse power $P_{X^+}$ in terms of such angle of rotation: a $\pi$-pulse thus corresponds to the first maximum of the resonance fluorescence emission, corresponding to a complete transfer from the hole to the trion state. Single photons can be emitted only if the trion transition is optically active and thus, if the QD is occupied by a single hole. Therefore, the detection of a single photon ensures that a single hole is confined in the QD.

Fig. 4(a) represents the time trace of the cross-polarized resonance fluorescence intensity emitted by one QD-pillar device, observed with a quasi-resonant power of $P_{QR} = 50 \mu \text{W}$, a $\pi$-pulse resonant excitation and a time bin of $\Delta t = 4 \mu \text{s}$. To clearly evidence the on/off behavior of this time trace, a second histogram representing the distribution of the number of detected events per time bin $\Delta t$ is represented in Fig. 4(b). When the hole is absent, photons are not emitted by the QD: this explains the high probability to detect zero photons per time bin $(7.3 \times 10^5$ events, in a total integration time of approximately 10s).

When the hole is trapped, the QD emits photons that are detected with a typical Poisson distribution centered on $\langle N \rangle = 7.9$ detected photons per time bin $\Delta t$, with a width $\sigma = 2.8(\approx \sqrt{\langle N \rangle})$. A small deviation from these two distributions is observed for $N = 1$, i.e. events where a single photon is detected during the time bin. This can be explained by the dark counts of the detector and by the imperfect polarization laser rejection. It is therefore possible to observe in real time the hole occupation by monitoring the trion resonance fluorescence signal.

It is in fact possible to determine the hole occupation probability and the hole trapping time with much better accuracy by measuring the second order correlations of the cross-polarized resonance fluorescence. A typical auto-correlation measurement is displayed in Fig. 5(a), where the two photon coincidence histogram is plotted for different detection delays. The peaks are separated by a delay $T_R = 1/f$, where $f = 82 \text{MHz}$ is the laser repetition rate. In Fig. 5(a), the continuous background between two peaks is due partly to the detector dark counts, but mainly to an imperfect filtering of the quasi-resonant laser.

Fig. 5(b) shows the same set of data for longer delays, at which we can observe that the envelope of the peaks is decaying. Fig. 5(c) shows the intensity auto-correlations $g^{(2)}(t)$ at even longer times (integrated with a time bin $\Delta t = 10 T_R$) with the same set of data (QD1 in red) and for two other devices (QD2 and QD3 respectively in purple and blue). Here, the large binning blurs out the peaks which were visible in Fig. 5(a) and (b), as well as the zero-delay antibunching. The correlation measurements evidence an exponential decay for the three devices, originating from the on/off fluctuations of the hole state in the QD. This effect is used to extract the hole confinement characteristics, as is shown in the following.

The auto-correlation function $g^{(2)}(t)$ can be interpreted as the probability to detect a photon in a detector (denoted APD1) at time $t$, conditioned to the detection at time $t = 0$ of a photon in the other detector (denoted QD)
ability that a hole is trapped in the QD is immediately after a photon detection in APD0, the probability occupation probability \( \langle P_h(t) \rangle = \frac{P_{h}(t)}{P_{APD1}} \) can be deduced from these correlation measurements.

APD0) normalized by the uncorrelated probability of detecting a photon at any time:

\[
g^{(2)}(t) = \frac{P(\text{APD}1, t| \text{APD}0, 0)}{P(\text{APD}1)}
\]

Immediately after a photon detection in APD0, the probability that a hole is trapped in the QD is \( P_h(0) = 1 \) and the detection of a photon in APD1 is thus more probable shortly after this first detection.

To interpret the dynamics of the charge state, we develop a theoretical model using two states: "0" (the crystall state, with zero excess charge) and "h" (the hole state). The CW quasi-resonant laser transfers the quantum dot state from "0" to "h" at a rate \( \gamma \) (which directly depends on the quasi-resonant power \( P_{QR} \), used to populate the single hole state). Conversely, the hole can tunnel out from the quantum dot with a tunneling time \( T_h \). The occupation probability of the empty state and the charged states are denoted \( P_0(t) \) and \( P_h(t) \) respectively, with \( P_0(t) + P_h(t) = 1 \). Their rate equations are:

\[
\begin{align*}
\frac{dP_h(t)}{dt} &= \gamma P_0(t) - \frac{1}{T_h} P_h(t) \\
\frac{dP_0(t)}{dt} &= -\gamma P_0(t) + \frac{1}{T_h} P_h(t)
\end{align*}
\]

We now assume that at time \( t = 0 \), a single photon emitted by the QD is detected. A hole is thus trapped inside the quantum dot with a probability \( P_h(0) = 1 \) and the hole occupation probability evolution is given by:

\[
P_h(t) = (1 - \langle P_h \rangle) e^{-t/\tau_{eff}} + \langle P_h \rangle
\]

with \( \langle P_h \rangle = P_h(\infty) = \gamma/\left(\gamma + \frac{1}{T_h}\right) \), the average hole occupation probability, and \( \tau_{eff} = \left(\gamma + \frac{1}{T_h}\right)^{-1} \), the effective time characterizing the charge dynamics.

In this model, the envelope of the auto-correlation function is given by:

\[
g^{(2)}(t) = \frac{P_h(t)}{\langle P_h \rangle} = \left( \frac{1}{\langle P_h \rangle} - 1 \right) e^{-t/\tau_{eff}} + 1
\]

Therefore, it is possible to extract the hole occupation probability and the hole tunneling time through these auto-correlation measurements, as illustrated in the inset of Fig. 5(c): the envelope of the \( g^{(2)}(t) \) varies from \( 1/\langle P_h \rangle \) at zero delay to 1 in a characteristic time \( \tau_{eff} \). In addition, its tangent at zero-delay cross the x-axis at \( t = T_h \).

It is also possible to take into account the small background noise and to obtain the real auto-correlation function of the quantum dot light source \( g^{(2)}(t) \) deduced from the experimental one \( g^{(2)}_{exp}(t) \). Let \( P_{QR} \) be the probability that a photon detected is originated from the quantum dot and \( 1-P_{QR} \), the probability that it is originated from bad laser filtering or dark counts. The relation between \( g^{(2)}_{exp}(t) \) and \( g^{(2)}(t) \) can be calculated and gives:

\[
g^{(2)}(t) = \frac{g^{(2)}_{exp}(t) - 2(1-P_{QR}) + (1-P_{QR})^2}{P_{QR}^2}
\]

In Fig. 5(c) and 6(a,b), the corrected \( g^{(2)}(t) \) is displayed, from which the occupation probability and the hole trapping time can be directly extracted using the fit with Eq. 4.

The dependence of these hole characteristics is now studied as a function of the quasi-resonant \( (P_{QR}) \) and resonant \( (P_{X+}) \) laser powers. Correlation measurements have been realized on QD1, for three different pulse areas of the resonant laser \( P_{X+} \) (\( \pi/2 \) and \( \pi/3 \) pulses), and with different quasi-resonant powers \( P_{QR} \). Fig. 6(a) shows the dependence of correlation measurements with respect to \( P_{QR} \), with \( P_{X+} \) fixed to \( \pi \)-pulse. Similarly, Fig. 6(b) shows the dependence of correlation measurements with \( P_{X+} \), with \( P_{QR} \) fixed to 50\( \mu \)W. These two
prorons, the setup transmission $\eta_{\text{det}}$.

The measured brightness should be linearly proportional to the hole occupation probability as a photon can be emitted only if a single hole is trapped inside the QD. This is indeed the case: the data corresponding to $\pi/2$, $\pi/3$ pulses are fitted altogether by 3 linear functions with relative slopes $s_{\pi/2} = s_{\pi/3} = s_{\pi} \cos^2(\pi/2)$ and $s_{\pi/3} = s_{\pi} \cos^2(\pi/3)$ (the $\cos^2$ takes into account the partial population inversion to the trion state for $\pi/2$ and $\pi/3$ pulses). For QD1, the maximum observed brightness is $21\%$ and corresponds to a $85 \pm 1\%$ occupation probability.

These experiments have been repeated on two other devices, QD2 and QD3. In these cases, a remarkably high polarized brightness has been measured: $B_p = 28 \pm 5\%$ for QD2 and $B_p = 33 \pm 5\%$ for QD3. They both show a similarly high hole occupation probability with $\langle P_h \rangle = 85 \pm 1\%$ for QD2 and with $\langle P_h \rangle = 91 \pm 1\%$ for QD3.

V. PERFORMANCES OPERATING AS SINGLE-PHOTON SOURCES.

Finally, the quantum properties of the cQED devices are investigated by assessing performances as single-photon sources. The single-photon purity is evaluated by the short timescale zero-delay intensity correlations, measuring the $g^{(2)}(0)$. A spectral filter (30pm bandwidth) is inserted in the collection setup to further suppress the spectrally-wide excitation laser and phonon sideband emission. The results obtained on QD1 show a good single-photon purity $g^{(2)}(0) = 1.6 \pm 0.4\%$ as displayed in Fig. 7(a).

The single-photon indistinguishability is evaluated by performing coalescence measurements via the Hong-Ou-Mandel effect. This is performed experimentally using
a path-unbalanced Mach-Zender interferometer\textsuperscript{53} where the difference of delay between the two arms is set to be equal to the laser period $T_R$. Therefore, two single-photons generated by two immediately separated laser pulses can interfere. Fig. 7(b) displays the experimental results obtained on QD1: the measured photon indistinguishability is also high, with an indistinguishability $V = 97 \pm 0.4\%$. These results illustrate the good quantum properties of the QD-photon interaction in the interfaces described in this paper.

VI. CONCLUSION

We have shown a technique for the deterministic realization of singly-charged QD-photon interfaces. The key points of our work are: facilitating the hole confinement thanks to a tunneling barrier; using in-plane magnetic-field spectroscopy to identify the trion transition prior to the etching step of the in-situ lithography process; and optically injecting a single hole with a quasi-resonant pumping scheme. The resulting spin-photon interfaces have then been used to monitor in real time the jumping of the hole in and out of the quantum dot. Thanks to autocorrelation measurements, we were able to obtain the hole tunnelling times in the tens of microseconds range, and the high hole occupation probabilities $\langle P_h \rangle$ between 85\% and 91\% for multiple samples. Moreover, we showed the good performance of bright sources of pure and indistinguishable photons, good indicators for high-quality spin-photon interfacing devices.

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