

OPTICAL PHYSICS

Toward optical quantum information processing with quantum dots coupled to microstructures [Invited]

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Received 22 February 2016; revised 29 May 2016; accepted 11 June 2016; posted 13 June 2016 (Doc. ID 259471); published 30 June 2016

Major improvements have been made on semiconductor quantum dot light sources recently and now they can be seen as a serious candidate for near-future scalable photonic quantum information processing experiments. The three key parameters of these photon sources for such applications have been pushed to extreme values: almost unity single-photon purity and photon indistinguishability, and high brightness. In this paper, we review the progress achieved recently on quantum-dot-based single-photon sources. We also review some quantum information experiments where entanglement processes are achieved using semiconductor quantum dots. © 2016 Optical Society of America

OCIS codes: (230.5590) Quantum-well, -wire and -dot devices; (270.0270) Quantum optics; (270.5585) Quantum information and processing; (270.5580) Quantum electrodynamics.

http://dx.doi.org/10.1364/JOSAB.33.00C160

1. INTRODUCTION

Photons are good candidates to perform optical quantum information processing as they can propagate over long distances with little dissipation. Their polarization can be viewed as a two-level system. Since a quantum bit (qubit) is a linear combination of basis states in a two-level system, written $\alpha|0\rangle + \beta|1\rangle$ following the Dirac notation with α and β complex numbers such that $\alpha^2 + \beta^2 = 1$, single-photon polarization can be used as a qubit. Moreover, only traditional polarizers and wave plates need to be used to initialize, manipulate, and project polarization-encoded single-qubit states.

The essential classical property for such single photons is the source brightness, while the essential quantum properties necessary for optically-based quantum information are singlephoton-*ness* of the light source and the identical-*ness* of all the single photons. To date, most quantum optics experiments, including those with quantum information applications, use sources based on heralded spontaneous parametric downconversion (SPDC), a spontaneous process that generates two photons into two spatial modes from one higher-energy photon. Typically, one of the downconverted photons is used to herald the other one in order to obtain nearly pure single-photon emission in the low pump power regime [1,2]. The indistinguishability of the photons is usually improved either by carefully engineering the phase matching of the SPDC process or by spectrally filtering the photons with a narrow bandpass filter to generate photon states with a coherence time much longer than the pump laser coherence length [3]. SPDC sources have been essential for many advanced photon-based quantum information experiments [4–6].

However, the low SPDC source brightness defined here as the probability to collect a photon per excitation pulse is 10^{-6} to 10⁻⁴, and is becoming a strong limiting factor for further quantum information experiments. This can be increased to 10^{-2} in continuous wave excitation using a heralded process [7]. The brightness can also be increased by increasing the pump power but the quantum properties of the photons are usually reduced at high brightness. The brightness is intrinsically limited by a multiphoton component that increases with the pump power [7]. For instance, boson sampling experiments on a linear circuit have been recently performed with up to 4-folds detections (four photons in four modes) but only few five-photon events have been detected, limiting this approach for larger systems [8]. One solution being explored is to improve the brightness at constant photon purity. It involves multiplexing several heralded photons from multiple SPDC sources into a single mode [7,9].

Developing other type of sources is an alternative approach. The conditions that such sources much satisfy are demanding: single-photon emission, indistinguishable photons, and bright emission. Nonheralded single-photon solid-state sources are one class of candidates. These includes organic molecules [10], nitrogen vacancy centers in diamond [11,12], or colloidal quantum dots [13,14] that can emit single photons at room temperatures.

This review focuses on a different solid-state photon source, self-assembled InAs/GaAs semiconductor quantum dots. They are a promising alternative light source for quantum information applications since they can have very pure single-photon emission [15] and they can emit highly indistinguishable photons [16]. They have high brightness values especially when coupled to microstructures. They can also emit entangled photons pairs [17,18]. In Section 2, we present the quantum and classical properties of these quantum dot sources and techniques to improve their brightness. Section 3 describes weak light-matter coupling (Purcell effect) using microcavities where enhanced quantum dot spontaneous emission is observed. In Section 4, we present the ultrabright sources of single and indistinguishable photons that have been developed in recent years. In Section 5, we describe the polarization entangled photons that can be emitted by a quantum dot. Finally, in Section 6, we review quantum optics experiments where quantum dots are at the heart of the process.

This review focuses specifically on the photons emitted from quantum dots and thus does not cover other important topics related to these quantum dots. The spin-photon physics associated with the charge excitons is one example [19–21], with a large body of on-going research that is beyond the scope of this review.

2. SEMICONDUCTOR QUANTUM DOTS FOR QUANTUM INFORMATION PROCESSING

Self-assembled InAs quantum dots are formed without lithographic patterning by strain-induced islanding, a process driven by the lattice mismatch between InAs and the host crystal, often GaAs [22]. InAs crystal growth initially occurs as a rough, planar region often called the wetting layer. As the accumulated strain increases with thickness island, growth becomes energetically favorable. Because of the lower bandgap of InAs with respect to the host crystal of GaAs, when the islands are overgrown, strong 3D quantum confinement results, forming nanoscale quantum dots.

A. Single Photons from InAs Quantum Dots

InAs-based quantum dots can emit single photons. The strong localization of the carriers inside the quantum dot enhances Coulomb interactions and the eigenenergies are strongly dependent on the carrier occupation in the quantum dot. Thus, the photon emission energies occurring during the carrier recombination processes depend on the quantum dot carrier occupation [23]. These quantum-confined states are excitonic states: a quasi-particle of localized Coulomb-bond electrons and holes. Examples include the single exciton (one electron–hole pair), biexciton (two electrons and two holes), or the singly charged exciton—the trion (one electron, one hole plus an extra electron or hole). Because of the 3D confinement and Pauli exclusion, the photons emitted from each of these excitonic states are single photons. The single-photon purity was demonstrated by Michler *et al.* in 2000 [15].

The single-photon purity is quantified through a normalized second-order autocorrelation measurement. It is traditionally measured using a Hanbury Brown and Twiss interferometer [24] where a beam splitter in the setup channels photons to

two detectors each with single-photon resolution. Secondorder correlations measurements, $g_{HBT}^{(2)}$, are made on photon detections. If τ is the time difference between photon detections, particularly important are the normalized second-order statistics at $\tau = 0$, $g_{\text{HBT}}^{(2)}[0]$. A value of $g_{\text{HBT}}^{(2)}[0] < 1$ is the hallmark of a quantum source. A value of zero uniquely identifies the state as a purely single-photon state. Nonzero values mean the photon state is a mixture of several photon number states, and values below 0.5 indicate a mixture of single-photon states [25]. Figure 1(a) shows an example of a $g_{\rm HBT}^{(2)}$ function obtained by exciting a quantum dot with a 12.2 ns repetition rate modelocked laser. Here the quantum dot is excited into a discrete state either a lower-energy state of the wetting layer or an excited state of the quantum dot. We call this quasi-resonant excitation. From the area of the correlation peak at delay 0, a $g_{\rm HBT}^{(2)}[0] = 0.01(1)$ is found indicating high single-photon purity of the source [27].

Further information on the photon statistics can be obtained by performing a two-time correlation function $g_{\rm HBT}^{(2)}(t_1, t_2)$ where t_1 and t_2 are the delay times between the excitation pulses and the photon detection on the two detectors. This function provides information on the dynamics of the photon emission. Using this technique on quantum dot photon emission, Flagg et al. [26] observed that when the excitation laser energy is larger than the quantum dot confined states that the single-photon purity depends on the time difference between the photon emission and the excitation laser pulse. Early emitted photons have a lower purity than the average [Fig. 1(b)]. This phenomenon is explained by delayed capture processes: even though the laser pump pulse (≈ 5 ps) is much shorter than the transition radiative lifetime, carriers generated in the GaAs or InAs wetting layer by the pump laser can decay through several pathways to the quantum dot. If one pathway results in early exciton emission the state can be repopulated through a slower pathway, and a second emission can occur before the next laser pulse. Values as small as $g^{(2)}_{\rm HBT}(0) = 3 \times 10^{-3}$ have been reported

Values as small as $g_{\rm HBT}^{(2)}(0) = 3 \times 10^{-3}$ have been reported recently using resonant fluorescence excitation [28–31], where recapture processes are eliminated. However, resonant excitation requires additional techniques to minimize the laser scattering, such as cross-polarized excitation and detection from above the sample [28,29,32,33] or side excitations where the light field is guided in a planar cavity made of two distributed Bragg reflectors [34–36].



Fig. 1. Single-photon characterization. (a) Example of a $g_{HBT}^{(2)}(\tau)$ autocorrelation function under a quasi-resonant pulsed excitation. (b) Two-time second-order correlation, $g_{HBT}^{(2)}(t_1, t_2)$, measured for an excitation in the wetting layer. Figure adapted from [26].

The single-photon purity required to perform quantum information processing depends on the protocol that is being used. For example, the polarization states of two photons can be entangled (so that the state of each photon cannot be described independently) using a polarizing beam splitter. The quality of the entanglement depends on the degree of indistinguishability of the photons (see Section 2.B) but also on the single-photon purity. One can show that such entangled states can violate Bell inequalities when $g_{\rm HBT}^{(2)} < 0.15$ [7,37]. The generation of a three-photons entangled state is usually more demanding and some protocols requires $g_{\rm HBT}^{(2)} < 0.06$ (for instance, to satisfy the witness value of a Green–Horne–Zeilinger (GHZ) state, i.e., a state generated with two polarizing beam splitters [38]). Some further advance protocols require even higher purities, such as $g_{\rm HBT}^{(2)} < 0.03$ for some deterministic two-qubits gates [7].

B. Indistinguishable Photons

High-purity single photons can be used for quantum information protocols, such as the BB84 protocol [39]. However, long-distance quantum communication systems will require quantum repeaters to beat the low, but nonnegligible, photon loss over long-distance fiber propagation. Such repeaters, installed along the communication channel, will distribute the entanglement over segments of the quantum channel [40,41]. In such cases, the photons must be indistinguishable. The photon indistinguishability is also required for many quantum information protocols, for instance, where qubits interactions need to occur to perform logic operations.

When two indistinguishable single photons simultaneously enter two separate ports of a 50:50 beam splitter, because photons are bosons, theory predicts the photons will coalesce into the same mode, thus exiting the beam splitter through the same output port [42,43]. The degree of indistinguishability, *C*, of two single photons is usually measured using a Hong–Ou– Mandel interferometer [44], an unbalanced interferometer where one path can be adjusted to adjust the time delay of photon interference at the output beam splitter. Phase stabilization is not required. Second-order time correlations, $g_{HOM}^{(2)}$, are performed between detected events at each output port. When the photons are indistinguishable and interfere at the beam splitter, no two-photon detection events occur. Thus, the value of $g_{HOM}^{(2)}[\tau = 0]$ depends on the indistinguishability, $g_{HOM}^{(2)}[0]$ goes to 0 for ideal quantum light and C goes to 1.

The indistinguishability of successively emitted quantum dot photons was shown by Santori *et al.* in 2002, only 2 years after the demonstration of the first quantum-dot-based single-photon source [16]. The experiment is shown in Fig. 2(a). A Hong–Ou–Mandel interferometer, as described above, and a pulsed excitation scheme are used. Several peaks appear in the $g_{HOM}^{(2)}[\tau]$ function because of the small, 2 ns delay and the different possible paths that can be taken by the photons. For lossless 50:50 beam splitters, the area of the peak at delay zero is proportional to $g_{HBT}^{(2)}[0]$ and it fully vanishes for an ideal quantum light.

The required photon indistinguishability necessary for quantum information processing experiments is ideally unity but the requirement can become relaxed in some cases. For instance, in the quantum controlled-not gate case presented in



Fig. 2. Measurement of indistinguishability. (a) Schematic of a setup to measure the indistinguishability of two photons emitted from a quantum dot that is excited 2 ns apart. They are sent to an unbalanced Michelson interferometer and then to detectors. (b) Resulting second-order correlation histogram of the quantum dot light. The intensity of the peak at delay $\tau = 0$ is proportional to the photon indistinguishability. Figures adapted from [16].

Section 6.B, an indistinguishability of 0.5 is the quantum limit for the generation of entangled photon pairs [Eq. (12) and Fig. 11].

Maximum quantum interference can occur only when the photons are spatially and temporarily overlapped, and have the same polarization. In addition, the photons must be Fourier transformed limited and have the same energy. The first two parameters depend on optical alignment while the last parameters are dependent on the quantum dot source.

Fourier transformed limited photons have a coherence time, T_2 , limited by their radiative lifetime, T_1 , such that $T_2 = 2T_1$. Any pure dephasings, with a characteristic time T_2^* , will decrease the coherence time and thus coherently broaden the transition linewidth, γ :

$$\gamma = \frac{1}{T_2} = \frac{1}{2T_1} + \frac{1}{T_2^*}.$$
 (1)

Pure dephasing effects occur at time scales shorter than the transition lifetime. On the contrary, slow decoherences lead to transition fluctuations and induce inhomogeneous broadening of the transition, Γ_{in} . These decoherences can result from charge fluctuations or nuclear spin flips in the quantum dot surrounding, leading to Zeeman or DC Stark shifts [45], which in turn reduces indistinguishability. A measurement integrated over several minutes will give access to a lower bound for the degree of indistinguishability of photons emitted only few nanoseconds apart [16].

At least three techniques can be used to improve the indistinguishability of quantum dot photons: reduce the radiative lifetime, use resonant fluorescence, and control charge fluctuations. In 2002, Santori *et al.* coupled a quantum dot to a microcavity in order to benefit from the Purcell enhancement to reduce the lifetime, T_1 [16]. Indeed, Eq. (1) shows that shorter lifetimes reduce the inhomogeneous spectral broadening induced by the pure dephasing T_2^* . They obtained indistinguishability values up to 81%. The temporal difference between the photons was small, 2 ns, about 10 times the radiative lifetime. Further cavity quantum electrodynamics advantages will be detailed in Sections 3 and 4. Resonant pumping schemes improved the indistinguishability to 0.97 (2) in 2013 [28]. Combining cavity quantum electrodynamics and resonant pumping has resulted in near unity indistinguishability (0.985) and high brightness [46]. A better control of the charge fluctuations in the quantum dot vicinity using electrically controlled micropillars [47] further increased the photons indistinguishability to 0.9956(45) with an extraction efficiency of 0.65[31]. Because somewhat similar values can be obtained without electrical bias, it remains to be seen if the bias is necessary in all cases to achieve such a high level of indistinguishability.

C. Brightness

Single and indistinguishable photons are two quantum properties of light that are necessary to ensure optimum quantum information processing operations. The brightness of a light source can be defined as the probability, p_1 , of collecting a single photon in the first external optic every laser pump pulse. A bright source is crucial to perform quantum operations on large computational Hilbert space with many fold correlations. The probability of distributing N photons into N different modes is, at best, $p_N^{(N)} = (p_1)^N$, and thus having a large p_1 is critical to any scaling. Moreover, a bright photon source is required for many advanced quantum information protocols like for the loophole-free Einstein–Podolsky–Rosen (EPR) experiment with entangled photons. The overall efficiency (source brightness, optics, and detection) needs to be above 82.8% in the EPR traditional experimental scheme [48].

The brightness, p_1 , and the extraction efficiency, η , are linked by

$$p_1 = p_s \times \eta, \tag{2}$$

where p_s is the probability of the quantum dot to be in a target state at every laser pulse. For example, the brightness of an exciton state will be reduced if the quantum dot is occasionally in a trion state, even if η is large. The term p_s will also be reduced in resonant pumping if the state spectrally wanders outside the laser linewidth since the probability of creating the target state resonantly will be reduced. There is no general method to increase low values of p_s but some tricks can be used, such as applying an external electric field or a weak above band laser, to partially control the quantum dot charge state and to restore its brightness [49–51].

InAs quantum dots are embedded in a GaAs medium with a refractive index of about 3.5. Thus, one can only expect an extraction efficiency of $\eta \approx 2\%$ (collection from the top surface) for quantum dots in bulk GaAs materials [52]. The growth of AlAs/GaAs distributed Bragg reflectors (DBRs) around the quantum dot to form a λ planar cavity (λ/n thick GaAs layer, *n* being the GaAs refractive index) can improve the collection efficiency to $\eta \approx 10\%$ –20% when the bottom DBR has a much higher reflectivity than the top one [52]

[Fig. 3(a)]. The collection efficiency can be improved, up to 10%-30%, over a wide frequency range by using solid immersion lenses [55,56].

When microstructures are designed around the quantum dot to further improve its brightness, the extraction efficiency, η , can be written:

$$\eta = \beta \times (1 - \alpha) \quad \text{with } \beta = \frac{\Gamma}{\Gamma + \Gamma_{\text{other}}}.$$
 (3)

The term $(1 - \alpha)$ is the fraction of the microstructure optical field that can be collected by the first lens of the collection setup. The term α includes scattering losses and the fraction of the field that cannot be collected with traditional optical setups. The β factor characterizes the fraction of the light emitted by the quantum dot that is coupled into the target microcavity mode. Γ and Γ_{other} are the spontaneous emission rates of the QD transition, respectively, into the microstructure mode and into all the other modes. Thus, increasing Γ or reducing Γ_{other} improves the β factor and hence the source brightness.

Increasing β through inhibition of Γ_{other} . The term Γ_{other} can be decreased by shaping the electromagnetic field around the quantum dot to inhibit the spontaneous emission. Inhibition effects have been observed in photonic crystal cavities [57] and in photonic crystal waveguides [58]. In addition, structures using a photonic crystal waveguide exhibit almost unity emission rates into the waveguide, up to $\beta = 0.98$ [59–61]. The outcoupling efficiencies to free space are not specified but quantum information processing could be performed directly using photonic crystal waveguides [62].

Structures using nanowires or inverted nanotrumpets can also inhibit the spontaneous emission of a quantum dot and increase the β factor [53,63,64] [Fig. 3(b)]. The inhibition is produced by reducing the nanowire diameter at the position of the quantum dot layer [65]. Brightness values up to $p_1 \approx 0.75$ have been measured [53]. These structures are broadband in frequency and thus, they could allow for entangled photon-pairs generation



Fig. 3. Extracting quantum dot photons. (a) A quantum dot sandwiched between two asymmetric DBR mirrors allows for up to \approx 10%–20% collection efficiency. (b) Brightness of \approx 75% has been seen with a quantum dot inserted in the bottom of an inverted trumpet structure. Figure adapted from [53]. (c) Ultrabright sources ($p_1 \approx$ 79%) of single and indistinguishability photons have been made by coupling a quantum dot to a \approx 3 µm micropillar cavity. Figure adapted from [54].

using the exciton and the biexciton photons [18,29,66,67]. However, the small distance between the quantum dot and the surface is a drawback that may degrade the photons indistinguishability. Some surface passivation methods have been developed to minimize drifts of the quantum dot states [68].

Increasing β **through the enhancement of** Γ . The spontaneous emission rate into a particular mode, Γ , can be enhanced by coupling the transition to a microcavity to benefit from the Purcell factor, as discussed in Section 3.A. Interestingly, this improves both the brightness and the photon indistinguishability. Extraction efficiencies of ≈ 0.5 have been measured with suspended circular Bragg grating microcavities [56,69] or photonic crystal cavities [70] and up to 0.79 using micropillar cavities (Section 4 and [54]).

Inhibition and enhancement. Some structures like confined Tamm plasmon modes exhibit both spontaneous emission enhancement of Γ and inhibition of Γ_{other} : the β factor can approach unity [71]. However, the term $(1 - \alpha) \approx 0.72$ can be a limiting factor of those structures [72].

3. ENHANCEMENT OF THE SPONTANEOUS EMISSION RATES WITH LIGHT-MATTER INTERACTIONS

Enhancement of spontaneous emission can be achieved by coupling a quantum dot to a microcavity because of the Purcell effect (Section 3.A). Several kinds of cavities have been developed and besides weak cavity-quantum dot coupling with enhanced spontaneous emission (Purcell effect), strong coupling has been observed (Section 3.B) using a number of techniques (Section 3.C).

A. Weak Coupling and the Purcell Factor

Weak light–matter interactions modify the spontaneous emission rate of an emitter compared to the free-space regime. Weak coupling occurs when the cavity losses are large enough so the coupling is only from the emitter to the cavity mode and is not reversible: $4\Omega < \omega/Q$ where Ω is the Rabi frequency, ω the emitter angular frequency, and Q the cavity quality factor [73]. If the emitter dipole is weakly coupled to a microcavity, its spontaneous emission rate can be enhanced by the Purcell effect [74] and the spontaneous emission rate $\Gamma = 1/T_1$ can be approximated by the Fermi Golden rule [74–76]:

$$\Gamma = \frac{2\pi}{\hbar^2} \langle |\langle \vec{d} \cdot \vec{E}(\vec{r}_{\rm QD}) \rangle|^2 \rangle \times \rho(\omega_{\rm QD}),$$
(4)

where $\rho(\omega_{\rm QD})$ is the density of photon modes at the emitter's frequency, $\omega_{\rm QD}$; the term \vec{E} is the electric field operator; and $\vec{r}_{\rm QD}$ is the location of the quantum dot dipole, \vec{d} . The outer averaging is performed over the modes seen by the emitter. A microcavity can play a key role here because it locally increases the density of electromagnetic states. The maximum of spontaneous enhancement is spatially located at the maximum of the electric field, $\vec{r_c}$, and at the cavity resonance frequency, ω_c . The maximum of spontaneous enhancement occurs when the quantum dot transition and the cavity mode are overlapped ($\vec{r}_{\rm QD} = \vec{r_c}$) and have the same frequency ($\omega_{\rm QD} = \omega_c$). In this case, the ratio of the spontaneous emission rate in the cavity mode, Γ , over to the coupling to free space, Γ_0 is proportional to the Purcell Factor F_p :

$$\frac{\Gamma}{\Gamma_0} \propto F_p = \frac{3Q(\lambda_c/n_{\rm eff})^3}{4\pi^2 V_{\rm eff}},$$
(5)

where λ_c is the wavelength of the cavity and n_{eff} the effective refractive index of the cavity. Therefore, cavities with high quality factors, Q, and small effective mode volumes, V_{eff} , can result in regimes with strong Purcell factors.

B. Microcavities

Several kinds of photonic microcavities have been coupled to semiconductor quantum dots. The most studied structures are micropillar cavities, photonic crystal cavities, and microdisk cavities. Many interesting properties and effects have been reported with these structures; in particular, radiative lifetime enhancement in the weak coupling regime [77-79] and normalmode splitting in the strong coupling regime between a single quantum dot state and the cavity [80-82]. More recently, new structures have been developed. For example, fiber-based external-mirror microcavities where the quantum dot epitaxial layer sits on a bottom DBR and a curved DBR is fabricated on a single-mode fiber [83]. The advantages are that a large fraction of the light is directly coupled into a single-mode fiber and, more notably, the cavity can be spatially and spectrally coupled to a specific quantum dot by moving the fiber. Suspended circular Bragg grating microcavities have also been developed and they exhibit good properties with the observation of a Purcell enhancement and extraction efficiencies of $\eta \approx 0.48$ [56,69].

Currently, micropillar cavity structures produce the highestquality light sources in terms of the quantum properties of single-photon purity and indistinguishability, and of brightness. A micropillar is a cylinder a few micrometers in diameter. Two DBR structures are located along the pillar axis to form a λ -cavity between them. The quantum dot layer is grown in the middle of this cavity, where the electric filed is maximum [Fig. 3(c)]. The longitudinal field confinement is created by the DBRs and it can be calculated using standard transfer-matrix calculations [84]. The transverse field is determined using similar methods as used for standard waveguides and the fundamental mode is a HE11 mode [76,85,86]. Numerical apertures are rather small, typically ≈ 0.4 for $\approx 3 \ \mu m$ diameter micropillars [87]. Coupling rates into a single-mode fiber higher than $\approx 70\%$ have been obtained by optimizing the mode matching [27].

C. Coupling of a Quantum Dot into a Microcavity Mode

The deterministic coupling of a single emitter to a micrometer structure requires control of the energy and the position of the emitter and of the cavity resonance (Section 3.A). Self-assembled semiconductor quantum dots, described above, are unfortunately randomly distributed on the growth layer surface. Moreover, their exciton energies are inhomogeneously distributed as they depend on the quantum dot size, shape, and composition that cannot be precisely controlled during the growth process [80,81,88]. Thus one must fabricate hundreds of cavities to hopefully get a few cavities coupled with a quantum dot. This method clearly suffers from the random spectral and spatial overlap between the quantum dots and the cavity modes.

High-resolution microscopy techniques. Scanning electron microscopy [89] and atomic force microscopy [90] techniques have been developed to precisely locate a quantum dot (\approx 30 nm accuracy) and to define a suspended photonic crystal cavity centered on the quantum dots [Fig. 4(a)]. Spectral matching is achieved by enlarging the photonic crystal holes and thinning the membrane using sequential etching processes.

Optical in situ lithography technique. Optical lithography at liquid helium temperatures can be used to define a mask on a sample surface. A micropillar cavity can then be etched so the cavity is spatially and spectrally aligned to a quantum dot [91]. The method uses a photoresist deposit on the sample surface and two lasers [Fig. 4(b)]. One laser, with a wavelength above ≈ 800 nm, can excite a quantum dot, but barely exposes the resist. The quantum dot photoluminescence signal is sent to a spectrometer and to a CCD camera. Looking at the quantum dot photon intensity as a function of the quantum dot's position relative to the laser, one can determine the emitter position within a ≈ 20 nm accuracy. The second laser, green at 532 nm, is overlapped on the first one and is used to expose the photoresist at the position of the quantum dot. Finally, coarse spectral matching is achieved by adjusting the photoresist exposed area since the cavity mode energy depends on the diameter of the micropillar [76,86] and is fine tuned by changing the sample temperature by a few degrees or the electric field applied to the quantum dot [82,92].

Photoluminescence imaging. More recently, another optical technique has been reported [56]. The main difference between this method with the optical *in situ* lithography technique is that here the optical processes are used to find a quantum dot of the desired wavelength range and its position is



Fig. 4. Coupling of a quantum dot to a microcavity. (a) Atomicforce microscope topography of a photonic crystal nanocavity aligned to a quantum dot. The small hill in the middle arises from a quantum dot (63 nm below the surface). The color bar indicates the measured height. Figure adapted from [90]. (b) Schematic of an optical in situ lithography technique. Two lasers are used to find the quantum dot position and energy, and to define a cavity around it. Figure adapted from [91]. (c) Image of the photoluminescence signal of a quantum dot centered with a circular Bragg grating microcavity. The scale bar represents 5 μ m. Figure adapted from [56].

noted relatively to a mark on the surface. A red laser and an infrared LED illuminate the sample surface and excite the quantum dot sample so that both the quantum dot light and the markers are visualized on a camera. The cavity, a circular Bragg grating microcavity, is then fabricated using the marker position and electron-beam lithography [Fig. 4(c)]. An extraction efficiency of $\eta = 0.40$ is reported.

A new technique was reported in 2015 [93]. A quantum dot is located using a low temperature cathodoluminescent technique [94]. A microlens is fabricated, *in situ*, using electronbeam lithography process. An indistinguishability value of 0.80 is reported and a photon extraction efficiency of $\eta = 0.23$.

We note that quantum dots can also be grown on predetermined sites [95,96]. The quantum properties of these quantum dots are still inferior to the self-assembled quantum dots but large improvements have been made [97]. More complicated integrated structures will likely require ordered quantum dot arrays and improvements in results from these techniques will be important.

4. BRIGHT SOURCES OF INDISTINGUISHABLE SINGLE PHOTONS USING MICROPILLAR CAVITIES

In Section 2, we have seen that semiconductor quantum dots can emit single photons (2000, [15]), indistinguishable photons (2002, [16]), and be a bright light source (2010, [63]). However, combining those three parameters remained a challenge. In this section, we summarize a method based on a micropillar structure (Section 3.B) that allows the combination of high brightness (Sections 4.A and 4.B) and high indistinguishability (Section 4.B) [54]. The quantum dot is deterministically coupled to the cavity by an optical *in situ* method using a technique described in Section 3.C. Further improvements of the quantum properties of the source were made by exciting the quantum dot resonantly during subsequent measurements [28,30,31,46].

A. Extracting Single Photons with a Micropillar Cavity

In the general case described in Section 2.C, the extraction efficiency η is equal to $\beta \times (1 - \alpha)$ with β being the fraction of the photons emitted into the cavity mode and $(1 - \alpha)$ the fraction of the cavity field that can be collected. In the case of a quantum dot coupled on resonance and at the maximum electric field position of a micropillar cavity, the Purcell factor, $F_P = \Gamma/\Gamma_0 \propto Q/V_{\text{eff}}$, where Q is the quality factor of the micropillar cavity and V_{eff} the effective volume of the cavity [Eq. (5)]. In this case $\Gamma_{\text{other}} = \Gamma_0$ because the emission into the other modes is barely modified by the cavity. Thus, from Eq. (3):

$$\beta = \frac{F_P}{F_P + 1}.$$
 (6)

In the case that the bottom DBR reflectivity is much larger than of the top DBR, most of the light is emitted through the top surface. The cavity losses are mostly due to scattering on the micropillar sidewalls and by light escaping the cavity through the top mirror. Losses induced by the DBR layers can be small because quality factors larger than 10^5 have been obtained with similar structures with a wider diameter [98]. Because losses reduce a cavity quality factor and broaden its linewidth, and because the quality factor of a planar cavity, Q_0 —an infinitely large micropillar—is mostly given by the losses on the mirror, we can directly find the fraction of the light collected by the first lens, $(1 - \alpha)$. It is the ratio of the quality factor in the confined cavity case, Q, and in planar cavity cases, Q_0 [54]:

$$(1-\alpha) = \frac{Q}{Q_0}.$$
 (7)

A fit to experimental measurements of Q/Q_0 is shown in Fig. 5. The curve drops for small diameters because the cavity mode field is laterally less confined and scattering due to the sidewall roughness is enhanced. In addition, when the diameter decreases to 1 µm, the β factor increases because the Purcell factor inversely depends on the effective cavity mode volume. The range of diameters [2–3] µm is a good compromise between losses by scattering and low Purcell factor. Theoretical efficiencies are ≈80% with such micropillar cavities [54,67]. Adiabatic micropillar cavities, where the DBR thickness is gradually changed around the quantum dot layer, should exhibit even better extraction efficiencies as the field is laterally better confined [99].

B. Bright Single-Photon Sources

Using the optical *in situ* lithography technique described in Section 3.C, micropillars with diameters around 3 μ m with a quantum dot coupled to the cavity were fabricated to optimize the extraction efficiency [54]. For the cavity discussed here, and reported in [54], the measured radiative lifetime was $T_1 = 265$ (30) ps [moderate Purcell factor, $F_p = 3.9$ (6)], a cavity quality factor $Q_0 \approx 3000$, and $Q/Q_0 = 0.95(5)$. During the lithography process the authors selected quantum dots to ensure $p_s > 0.95$. Thus, a brightness value of $p_1 = 0.75$ (16) is expected.

The brightness is experimentally measured by exciting the sample using a pulsed laser at 860 nm (\approx 3 ps at a rate of 82 MHz). The photons are collected with a microscope objective, sent to a spectrometer, and to a single-photon avalanche photodiode. The brightness of the source, p_1 —the probability



Fig. 5. Brightness optimization. Fit of the experimentally measured $(1 - \alpha) = Q/Q_0$ terms (black dashed line), calculated $\beta = F_P/(F_P + 1)$ (red dotted line), and the maximum theoretical extraction efficiency $\beta \times (1 - \alpha)$ (solid green line) as a function of micropillar diameter. Figure adapted from [54].

to collect a photon per pulse at the first lens of the microscope objective—is obtained by normalizing the count rate measured on the detector by the laser repetition rate (82 MHz) and the setup detection efficiency (typically $\approx 1\%$). The maximum brightness of the source, p_1^{max} , is obtained in the saturation regime. In this regime, a photon is emitted from a target state whose occupation probability is maximum (Section 2.C). In lower pump regimes, the occupation probability drops and the brightness equals

$$p_1 = p_1^{\max}(1 - e^{-P/P_{\text{sat}}}),$$
 (8)

where *P* is the laser power and *P*_{sat} the saturation power. The experimental data are plotted in Fig. 6(a) (dotted line). The single-photon purity is measured simultaneously to account for multicapture processes within the same laser pulse [26]. For this they used the correction coefficient $\sqrt{1-g_{\rm HBT}^{(2)}[\tau=0]}$ [100]. Figure 6(b) shows $g_{\rm HBT}^{(2)}[\tau=0]$ as a function of the pump power. Values are always smaller than 0.15 so the correction factor is small. Finally, the maximum brightness value measured from this source was 0.78 (8) [Fig. 6(a), solid line], the highest brightness reported to date. More recent works reported similar values of about 0.65 [31,46].

C. Historical Development of Bright Indistinguishable Photons Sources

Combining high brightness and high indistinguishability is not trivial. As visible in Eq. (8), to obtain maximum brightness, the pump power must be adequately strong so the system is in the saturation regime. However, this strong pump power creates carriers that may dynamically alter the local potential landscape around the quantum dot and decrease the photon indistinguishability.



Fig. 6. Current brightness results. (Top) Raw (open squares) and multiphoton corrected (solid squares) number of collected photons per laser pulse and the corresponding detected count rate per second. (Bottom) Values of $g_{HBT}^{(2)}[0]$ values as a function of the pump power. The excitation laser is tuned to 860 nm to create carriers in the wetting layer. Figure adapted from [54].

In Fig. 7(a), we summarize some of the main achievements toward bright sources of indistinguishable photons. Some results have improved the photons indistinguishability [16,28,30,103,104], and others have improved the brightness [53,63,101,102], but none simultaneously. Simultaneous improvements of the two properties started to be reported in 2013 in [54] from the Senellart group where several techniques were explored to combine high brightness and high photon indistinguishability. We review some of them here.

Pumping into the wetting layer [105,106] was tried and the results of the indistinguishability measurements are plotted in Fig. 7(b) as a function of the pump power. Values of the indistinguishability of successively emitted photons as high as 0.86 (10) are shown at low pump powers, but the indistinguishability drops down to ≈ 0.50 when the brightness of the source is increased. One explanation is because of the high laser power required to reach the saturation regime [right scale in Fig. 7(b)]. Thus, excess carriers can be generated in the quantum dot surrounding leading to spectral diffusions and decoherences.

A solution to avoid the excessive generation of carriers is to excite the sample using a quasi-resonant pumping scheme. Several groups have reported indistinguishability values between 0.7 and 0.8 under such pumping regimes [16,103,104]. However, in [54], indistinguishability values of about 0.5 were reported at almost any excitation powers [triangles in Fig. 7(b)]. Likely, this is due to sample variation for this technique.

Inspired by the weak above-band laser technique used to restore the source brightness under resonant excitation [50,51], a two-color excitation scheme with a weak wetting layer wavelength laser (860 nm) and a quasi-resonant laser (906 nm) was used in [54]. This technique helped to improve the indistinguishability of the photons making this source the first single-photon source to be both bright and emit highly indistinguishable photons [stars in Fig. 7(b)]. High indistinguishability values up to 0.92 (11) were obtained and are almost a factor of 2 higher than without the weak 860 nm laser. The indistinguishability remains large, 0.82 (11) at higher source brightness, $p_1 = 0.65$ (6) photons collected per pulse.

D. Current State of the Art: Bright Indistinguishable Sources

Almost three years later, further improvements were made by several groups and recent works report almost unity photon indistinguishability at high source brightness [31,46,107]. Resonant excitation schemes and micropillar cavities are used.

Eliminating the laser scattering into detectors is challenging but a cross-polarization scheme in the excitation and collection paths, along with narrowband filters have produced large extinction ratios (>10⁵) between the pump laser and collected emission. Single-photon purity as low as $g_{HBT}^{(2)}[0] = 0.0028$ (12) has been reported [31]. An indistinguishability value for successively emitted photons of 0.985 at an extraction efficiency of 66% (the brightness is not specified) has been recently reported [46] with a quantum dot coupled to a micropillar cavity. An even higher indistinguishability value of 0.9956 (45) at a measured brightness of 0.154 (15) (extraction efficiency of 65%) was reported in late 2015 using an electrically controlled device to further minimize the charge fluctuations [31] [see historical plot, Fig. 7(a), and also Fig. 8].

5. ENTANGLED PHOTON SOURCES

Many photonics-based quantum information protocols require entangled states of light. For example, in long-distance quantum communication the single-photon loss can be overcome with quantum repeaters and these quantum repeaters require entangled states [40]. Several quantum computing protocols require entangled photons [108,109]; they are also required in quantum teleportation [1,110] and entanglement swapping experiments [111,112].



Fig. 7. Progress toward high brightness and high indistinguishability single-photon sources. (a) Blue points: indistinguishability of successively emitted photons and extraction efficiency of some quantum-dot-based sources reported since the first single-photon demonstration in 2000 [15] (yellow point). In 2002, the indistinguishability of the photons was reported [16]. Corresponding references: 2002-2007: [101,102]; 2010-2013: [53,63]; 2013: [54]; 2013-2014: [28,30]; 2015: [31,46]. The shaded area indicates the use of resonant fluorescence excitations. The red star is for good SPDC sources with $g_{\rm HBT}^{(2)} \approx 0.1$ [7]. (b) Indistinguishability values as a function of the brightness for different pumping conditions. Green squares, wetting layer pumping; red triangles, quasi-resonant pumping; blue stars, two-color scheme (see text). The solid black line plots the normalized laser power $P/P_{\rm sat}$ as a function of the source brightness. Figure adapted from [54].



Fig. 8. Almost unity indistinguishability; resonant fluorescence. (a) Schematic of a structure that emits highly indistinguishable photon under resonant excitation and electrical control. (b),(c) Second-order correlation histograms, $g_{HOM}^{(2)}$. The photon indistinguishability is determined by comparing the zero delay peak amplitude when the polarization in the two interferometer arms are parallel (b) and orthogonal (c) cases. In the parallel case, the zero delay peak should vanish for fully indistinguishable photons. Figures adapted from [31].

Entangled photon states occur naturally in the parametric downconversion process [113], and in an ideal quantum dot cascade from the biexciton state, through the exciton state and to the ground state. Decay from the biexciton state, $|XX\rangle$ occurs radiatively through two possible channels via the two neutral exciton states, $|X\rangle$. In an ideal, perfectly cylindrical quantum dot these $|X\rangle$ states are energy degenerate, but in practice are energy split by shape anisotropy (the anisotropic exchange splitting). If the frequency of this splitting is small compared to the radiative decay rate, information about the actual decay path is not available, and polarization entanglement can result [114].

The photon state generated from the quantum dot $|XX\rangle \rightarrow |X\rangle \rightarrow |0\rangle$ decay is a maximally entangled Bell state in the polarization basis { $|H\rangle$, $|V\rangle$ }:

$$\frac{1}{\sqrt{2}}(|H_{XX}H_X\rangle + |V_{XX}V_X\rangle). \tag{9}$$

If the anisotropic exchange splitting is larger than the transition linewidth the degree of entanglement is reduced. A timedependent dephasing term can be introduced between the terms $|H_{XX}H_X\rangle$ and $|V_{XX}V_X\rangle$ in Eq. (9) to account for this [115].

The first experimental demonstration of this entanglement was in 2006 [66], followed quickly by a postselection experiment [18]. Muller *et al.* showed how to eliminate the anisotropic exchange splitting by using a far-detuned dressing laser in the AC Stark regime [35], while other researchers used strain and electric fields to remove the splitting [116].

Cavity-enhanced entangled photon states are difficult to construct because the cavity must be resonant with both the $|XX\rangle$ and $|X\rangle$ states. This has been accomplished with a unique coupled pillar cavity in 2010 [67]. Research continues, including resonant excitation directly into the $|XX\rangle$ state through a two-photon absorption process [29], which is discussed in more detail in the context of time-bin entanglement.

6. TOWARD QUANTUM INFORMATION PROCESSING APPLICATIONS

The properties we reviewed here suggest that quantum light sources based on semiconductor quantum dots are possible candidates for some types of optical quantum information processing. Much work remains in the development of these sources and drawbacks still need to be solved. Still, these sources have evolved to the point where they can be used in some quantum information applications. We show some examples here.

First, we describe a time-bin entanglement protocol that leads to the generation of two entangled photons from a single quantum dot source. It uses the $|XX\rangle \rightarrow |X\rangle \rightarrow |0\rangle$ cascade discussed above in Section 5, but creates entanglement in time and energy rather than polarization and energy [117]. Second, we review an early realization of an entangling controlled-not gate that can entangle quantum dot single photons [27], and of a quantum dot device that acts as an optical controlled-not gate [118].

A. Time-Bin Entangled Photon Pairs

Polarization entangled states can be used in quantum computing; however, polarization can degrade in some cases, for instance, in long-distance optical transmission. An alternative approach is time-bin entangled states. In particular, such states are robust against dephasings in optical fibers and can propagate over long distances [119,120]. Time-bin entangled photon sources are currently mainly based on parametric downconversion processes [121,122].

It was proposed theoretically in 2005 that coherent resonant excitation of a quantum dot state can be used to generate timebin entangled single-photon pairs from two laser pulses [123]. The generation of time-bin entangled states requires the photons to be emitted at well-defined times, within the time windows of the $|Early\rangle$ and $|Late\rangle$ states. One challenging task is to coherently generate the photon pairs so that the emission time and phase remain unknown.

Here we highlight experiments by the Weihs group [117]. The photon pairs are emitted from the sequential recombination of the quantum dot $|XX\rangle$ and $|X\rangle$ states, and only one polarization state is used [Fig. 9(a)]. An advantage of timebin entanglement over polarization entanglement is that it does not suffer from the anisotropic exchange splitting, *s*, of the $|X\rangle$ quantum dot level since only one of the two relaxation paths is needed.

Time binning is achieved by resonantly pumping the quantum dot $|XX\rangle$ state [124]. A two-photon absorption process is used via a virtual state at half the $|XX\rangle$ state energy [Fig. 9(a)] [125]. Two peaks corresponding to the X and XX photon energies are visible in a spectrum [Fig. 9(b)]. The quantum dot is located in the middle of a planar cavity made of two DBRs to enhance the collection efficiency. The pulsed resonant laser scattering is reduced by exciting the quantum dot laterally from the side via the in-plane planar cavity waveguide and by using a pulse shaper to adjust the laser spectral width (Fig. 9).

When the two time-bin states, |Early and |Late are created by two coherent laser pulses generated by an unbalanced interferometer in the pump path, the output state is



Fig. 9. (a). Energy-level scheme for resonant two-photon excitation of a quantum dot biexciton state, $|XX\rangle$. A pair of photons is emitted in cascade through the exciton state, $|X\rangle$, to the ground state, $|0\rangle$. The term *s* is the anisotropic exchange splitting of the $|X\rangle$ level. (b) Quantum dot emission spectrum under resonant excitation of the $|XX\rangle$ state. (c) Real part of the reconstructed density matrix. The imaginary part is plotted in [117]. (b),(c) Figures adapted from [117].

$$\Phi = \frac{1}{\sqrt{2}} (|\text{Early}\rangle_{XX} |\text{Early}\rangle_{X} + e^{i\phi_{p}} |\text{Late}\rangle_{XX} |\text{Late}\rangle_{X}), \quad (10)$$

where $\phi_p = E_{XX} \Delta t/\hbar$ is the phase in the pump interferometer $(E_{XX}$ is the biexciton photon energy and Δt the delay in the pump interferometer). This state is a maximally entangled Bell state when $\phi_p = 0$ or π .

The quality of the time-bin entangled states is characterized with a quantum-state tomography measurement to reconstruct the density matrix [126]. Two additional unbalanced interferometers are installed for each XX and X photons. The delays on the excitation and on the two analyze interferometers are set equal and the relative phases between the three interferometers have to be stabilized during the measurement. An active stabilization system can be used, but in [117] the three interferometers share the same optical elements and operate in different spatial modes. The state tomography is performed is three bases by adjusting the phase in the X and XX analyze interferometers. A fidelity of the generated state with respect to the Bell state, Φ^+ [Eq. (10) with $\phi_p = 0$], of 0.69 (3) was measured. The real part of the density matrix is plotted in Fig. 9(c). The authors attribute the reduced visibility to a nonnegligible double excitation (on the "Early" and "Late" events) and dephasings. Coupling to a 3D microcavity should improve the fidelity because a higher photon extraction efficiency would allow the system to operate in a lower excitation power regime reducing the double excitation events.

B. Controlled-not Gates

Sets of single- and two-qubit quantum gates can be combined to create any other quantum gates [127,128]. Single-qubit gates

are easy to implement in the polarization space since a polarizer and half- and quarter-wave plates can define any polarization states of the Poincaré sphere. Two-qubit gates are more complex to implement because they require quantum interference between single photons. Knill *et al.* proposed in 2001 that any optical multi-qubit gate, in particular the two-qubit gate, could be implemented using linear optic components, projective measurements, and feed-forwards [109].

A quantum controlled-not gate is the main two-qubit entangling gate. It acts on two qubits, one commonly called the target qubit and the other the controlled qubit. This gate flips the target qubit conditionally on the control qubit state. A simplified controlled-not gate scheme was proposed in 2002 [129] and realized experimentally in 2003 [4]. Quantum interferences between the control and target photons can occur for some control qubit states and lead to a conditional target state flip. Several other schemes have been proposed theoretically and realized experimentally [130–135]. All of these referenced experiments were performed using parametric downconversion sources.

1. Optical Entangling Controlled-not Gate

The schematic of the gate reviewed in this section was proposed in [4] and is seen in Fig. 10. The key elements are the two calcite crystals that displace one linear polarization and transmit the other. These crystals improve the temporal stability of the interferometers and they transform the polarization encoded input qubits into path encoded qubits. The half-wave plate installed between the two crystals simulates the 1/3:2/3 beam splitter that is the main element of this gate scheme [129].

The gate was operated using a quantum dot source of bright indistinguishable single photons, as discussed above and reported in 2013 [27,54]. Two photons successively emitted by the quantum dot are probabilistically separated using a fiber beam splitter and delayed to arrive simultaneously on the controlled-not gate input ports. The target and control qubit photons are individually prepared, in any polarization state, using polarizers and quarter- and half-wave plates. The output state is monitored using quarter- and half-wave plates to project and detect any states on polarizers and single-photon avalanche detectors. Second-order correlations are performed on the detector outputs. This gate is probabilistic and only works when detections occur at each control and target outputs. The result of the operations are read on a second-order correlation measurement at delay 0 [129]. The truth table of the gate is determined in the rectilinear polarization basis, $[|H\rangle, |V\rangle]$,



Fig. 10. Optical quantum controlled-not gate. Schematic of the gate that was used in [27]. The key element for the stability is the calcite crystal and the main element of the gate is the central half-wave plates (blue lines) [129]. Two half-wave plates (blue lines) at 45° off their optical axis in the input and output target ports act as Hadamard gates.

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by measuring the output state for the four possible input states: $\{|HH\rangle, |HV\rangle, |VH\rangle, |VV\rangle\}$ (where $|ct\rangle$ are the control and target qubit states). Overlaps between the measured and the ideal truth table up to 0.73 were observed.

Truth tables have also been measured in one or two bases in [28,136] using quantum dot photons as well, but the entangling capabilities of the gates have not been reported.

Entangling controlled-not gate. A quantum controllednot gate can entangle photons that were initially independent. Indeed, if the control qubit is in a diagonal state $\frac{1}{2}(|H\rangle + |V\rangle)$ and the target qubit in the state $|H\rangle$, then the ideal output state is a maximally entangled Bell state:

$$\Phi^+ = \frac{1}{2} (|HH\rangle + |VV\rangle). \tag{11}$$

To show entanglement of single photons that are initially independent using this controlled-not gate, quantum dot photons were sent into the gate and the fidelity of the output state compared to the Bell state Φ^+ was measured. The control photons are prepared in the state $\frac{1}{2}(|H\rangle + |V\rangle)$ and the target photon in the state $|H\rangle$. The fidelity of the generated state is determined by measuring the degree of the correlations in three bases (rectilinear $|H\rangle$, $|V\rangle$; diagonal $|H\rangle \pm |V\rangle$; and circular $|H\rangle \pm i|V\rangle$) [137,138]. Figure 11(a) plots the fidelity as a function of the postselected source brightness. The postselected brightness is defined as the product of the brightness, p_1 , and the fraction of postselected photons. The postselection was performed on the arrival time of the photons with the lower postselected brightness values corresponding to shorter time differences. Without postselection and at a source brightness as high as 0.5 photons collected per pulse, the fidelity referenced to the Bell state Φ^+ is above the 0.5 limit for quantum correlations. With temporal postselection, the fidelity increases as the photon indistinguishability is improved [26] and at a postselected brightness of $p_1 = 0.15$, the fidelity reaches 0.71.

Perspectives. The quantum dot source used for this experiment emits photons with an indistinguishability of about 0.71 with postselection and 0.5 without. Following [129] to calculate the output coincidences for all the polarization configurations



Fig. 11. Entangling controlled-not gate. (a) Measured fidelity of the generated state compared to the Bell state ϕ^+ as a function of the source postselected brightness. The dotted line indicates the quantum correlation threshold. (b) Calculated fidelity as a function of the photons indistinguishability. The square indicates a fully distinguishable single-photon source; the circle (triangle) corresponds to the two experimental points without (with) temporal postselection. Figures adapted from [27].

and as a function of the indistinguishably, C, one can obtain a fidelity referenced to the Bell state Φ^+ of [27]

$$F_{\Phi^+} = \frac{1+C}{2(2-C)}.$$
 (12)

This function is plotted in Fig. 11(b). Using the latest sources discussed in Section 4.B where indistinguishability values are above 0.99 [31,46], the fidelity should be extremely close to one.

2. Controlled-not Gate Using Quantum Dot-Cavity Strong Coupling

A final example of a quantum information device utilizing quantum dots is also a controlled-not gate. However, while the last example of a controlled-not gate used quantum dot photons as the input to the controlled-not gate, in the work of Waks *et al.* discussed here the quantum dot–cavity system acts as the gate [118].

There are several differences between the system used by [118] and the previous quantum-cavity system. Here the microcavity is a photonic-crystal cavity (PCC) formed in a GaAs slab, less than 200 nm thick and suspended by etching a sacrificial AlGaAs layer under the GaAs. The cavity is completed in the standard PCC way by etching a periodic array of holes into the GaAs to form the in-plane confinement. Researchers search for a cavity with an InAs quantum dot in the defect region (an area without a hole) of the structure. Unlike previous devices discussed, here the quantum dot-microcavity system is in the strong coupling regime, and the success of the device is based on the cooperativity, $C = 2 g^2 / \kappa \Gamma$, where g is the cavity– quantum dot coupling strength, κ and Γ are the cavity and the quantum dot radiative decay rates. In [118] $g/2\pi =$ 12.9 GHz, $\kappa/2\pi = 31.94$ GHz, and $\Gamma/2\pi = 5.2$ GHz. Since $g > \kappa/4$ the system is in the strong-coupling regime, with C = 2.0. A similar experiment in a micropillar cavity showed a similar cooperativity (2.5) but in the weak-coupling regime [139].

Here, the controlled-not gate relies on the change in reflectivity when one of the two neutral exciton states is on-resonance with the cavity and the other is significantly off-resonance with the cavity. With a magnetic field of 1.6 T oriented along the growth direction (often called the Faraday geometry), one of the exciton states is tuned on-resonance with the cavity. Because of the magnetic field, the exciton transitions have circular polarizations, σ_+ or σ_- . Input states are $|H\rangle$ and $|V\rangle$, rotated $\pi/4$ to cavity polarization axis, $|x\rangle$ and $|y\rangle$, so that $|H\rangle = \frac{1}{\sqrt{2}}(|x\rangle + |y\rangle)$ and $|V\rangle = \frac{1}{\sqrt{2}}(-|x\rangle + |y\rangle)$. The important observation is that upon reflection, the states are transformed to $|H\rangle = \frac{1}{\sqrt{2}}(r|x\rangle + |y\rangle)$ and $|V\rangle = \frac{1}{\sqrt{2}}(-r|x\rangle + |y\rangle)$, where r is the reflectivity.

The logic goes as follows for light incident on the cavity (Fig. 12). If the quantum dot is in the exciton state that is detuned from the cavity, no transitions at the cavity resonance can occur since the off-resonant exciton is occupied. However, if the quantum dot is in the ground state (the empty state), strong cavity absorption will take place, modifying the reflectivity. How much the reflectivity is modified depends on C since r = (C - 1)/(C + 1). If $C \gg 1$, $r \to 1$. Thus, in this ideal situation under one condition the incoming state is



Fig. 12. Controlled-not gate. When a photon is sent to a coupled quantum dot cavity system, the phase of the reflected photon depends on the quantum dot state. If the quantum dot is in the ground state $|g\rangle$ and the cooperativity, $C \gg 1$, no phase is added, r = 1 and the incoming state in unchanged (top). However, if the quantum dot is in the excited state $|-\rangle$, r = -1 and the incident state is flipped (bottom). Figure adapted from [118].

flipped (r = -1), while in the other condition the state is reflected unchanged (r = 1).

Results show that when the system is in the off-resonant state, the qubit-flip probabilities are $P_{H\rightarrow V} = 0.93$ (3) and $P_{V\rightarrow H} = 0.98$ (3). This is due to the bare cavity effect since the magnetic field shifts the σ_{-} state far off the cavity resonance. Conversely, when the system in the ground state and incident light puts the system in the strong-coupling regime, the reflected state should ideally be unchanged. Here the probabilities of returning the same state are lower, $P_{V\rightarrow V} = 0.58$ (4) and $P_{H\rightarrow H} = 0.61$ (7). This makes good sense since while C is very high for a quantum dot microcavity system, C = 2, it falls short of the $C \gg 1$ ideal condition, so that r = 1/3, not 1. As with the other two demonstrations, this is an excellent first step but with improvements to the cooperatively better results can be expected in the future.

7. CONCLUSION AND PERSPECTIVES

In this review paper, we described some major progress made on quantum-dot-based light sources. These sources are still under development. For instance, there is ongoing research to improve the quality of the indistinguishability of photons emitted from the source very far apart temporally. Recently, photon indistinguishability values above 0.85 have been measured with photons emitted 160 ns apart [140] and very recently extended to over 10 μ s with indistinguishability of about 0.92 [141]. Long temporal streams of more than five highly indistinguishable photons will allow new possibilities for quantum information processing. For instance, they could be sent to a multiport quantum circuit to perform scalable quantum information processing on a large Hilbert space. One could think of the boson sampling experiments that are currently limited by the source brightness [8,142,143–145].

The interference of photons coming from several disparate quantum-dot-based sources is a key element for long-distance quantum communications, such as optical quantum teleportation and entanglement swapping protocols [1,112]. Interferences between photons emitted from distant quantum dots have been performed since 2010 [146–149]. Interference visibilities are typically in the range 25%–40%, and go up to

about 80% with spectral filtering; however, this strongly reduces the source brightness of the single-photon emission [150]. The generation of cluster states is an alternative scheme of quantum computation and it might be possible to obtain large-scale cluster state by coupling several quantum dots [151–153].

Several recent experiments are trying to better understand the spectral jittering effects that are a main limitation for interfering photons emitted far apart temporally (Section 2.B). To that aim, the temporal dynamics of photon energies and linewidths are being studied [45,154].

Major progress has been made on the semiconductor quantum-dot-based sources since the first demonstration of their single-photon properties in 2000. The coupling with microcavities and the use of resonant fluorescence excitation strongly improved their properties and now, these sources can deterministically emit single photons with almost unity purity and indistinguishability, and at high brightness values. Quantumdot-based sources will be used more and more in optical quantum information processing experiments, experiments that were up to some years ago almost exclusively reserved for spontaneous parametric downconversion sources.

Funding. National Science Foundation (NSF) through PFC@JQI; Army Research Office (ARO) Multidisciplinary University Research Initiative on Hybrid quantum interactions.

Acknowledgment. We greatly appreciate the support of P. Senellart here. A portion of this review is based on work O. G. conducted in the research group of P. Senellart at the Laboratoire de Photonique et de Nanostructures, CNRS, Marcoussis, France.

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