

Chip-Integrated Nanowire Quantum Dot Emitters for Scalable Quantum Photonics

by

Edith Pui Cheung Yeung

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Department of Physics
Faculty of Science
University of Ottawa

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Author's Declaration

This thesis consists of material all of which I authored or co-authored: see Statement of Contributions included in the thesis.

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Abstract

Single-photons are a fundamental component in many quantum optical technologies. An example of where single-photons are needed is in quantum cryptographic technologies such as the BB84 quantum key distribution protocol. This involves encoding data in the polarization of single-photons to share a key between the two parties. Such applications require a single-photon source that ideally can emit exactly one photon on-demand with a well-defined polarization and spatiotemporal mode. Among the various candidates, semiconductor quantum dots are two-level systems that can be tailored to emit single-photons at desired wavelengths. However, since the quantum dots emit randomly in all directions, they are typically embedded into different photonic structures for more efficient coupling to optical fibre systems.

In particular, bottom-up III-V nanowires embedded with quantum dots are promising single-photon sources for integration with large-scale quantum photonic circuits. To test the feasibility of a nanowire-Si₃N₄ waveguide integration, simulations of the power transfer from the fundamental HE₁₁ mode of a tapered nanowire to an underlying ridge waveguide predicted $> 95\%$ coupling efficiency. The mode within the nanowire is expanded as the light propagates up the nanowire taper, making it compatible for both collection with a free-space objective or to evanescently couple to an in-plane straight waveguide. In this thesis, different hybrid quantum photonic integrated circuit devices are demonstrated to efficiently harvest single-photons generated on-chip.

Through a pick-and-place method, individual nanowires can be transferred from the growth substrate to a chip containing Si₃N₄ photonic waveguides. This hybrid integrated device combines the efficient emission of quantum dots with the low-loss characteristics of waveguides. Coupling light on and off the chip is performed using lensed fibres. An evanescent coupling efficiency of $\eta_c \sim 84.8\%$ was measured with one device. Indistinguishable photons were successfully generated on-chip using above-band excitation and quasi-resonant (p-shell) excitation to obtain two-photon visibilities of 9.2% and a much improved 19.2%, respectively, over the temporal extent of the emitted photons.

A major limitation of these devices is that single-photons are more efficiently collected from the tapered end of the nanowire as opposed to the base. To collect the photons emitted in both directions, a curved waveguide with a long untapered nanowire can be used. This requires the precise placement of the quantum dot adjacent to the curved waveguide. To increase the tolerance for the quantum dot position along the curve, a curved segmented waveguide was investigated. A coupling efficiency of 63% was obtained from the combined counts of both the X and X^- transitions in a single dot.

For large-scale quantum networks, stable single-photon sources with a small footprint are desirable. The need for active alignment is eliminated with a plug-and-play source which allows for easy access to multiple fibre-coupled sources simultaneously. A fibre-packaged chip consisting of straight waveguides and tapered nanowire-quantum dots is demonstrated to have a total plug-and-play efficiency of $3.6 \pm 0.7\%$ simultaneously collecting from both ends of one device. The multiphoton emission probability of the device is $g^{(2)}(0) = 7.2\%$. The fibre-packaged chip was robust against multiple cooldown cycles and showed stable alignment over 27 hours of continuous operation.

The devices demonstrated in this thesis serve as proof-of-concept platforms for in-plane single-photon generation. Based on this work, single photons can then be manipulated with standard photonic components including ring-resonators, Mach-Zehnder interferometers, phase shifters for more complex applications on-chip.

Statement of Contributions

All the research presented in this thesis was conducted at the Quantum Optics lab at the National Research Council of Canada. Facility access to laboratory equipment was provided by Dr. Robin L. Williams and Dr. Dan Dalacu. The single-photon sources used in this thesis (i.e., nanowire-quantum dot samples) were grown by Dr. Philip J. Poole. Various foundries were used to fabricate the photonic integrated circuits presented, including: AIM Photonics, LIGENTECH, and LioniX International.

This thesis was entirely written by the author, who also performed all data collection, analysis, and plotting. The author has attempted to acknowledge all contributions from the collaborators at the appropriate points in each chapter. The thesis contains the following papers:

1. **Yeung, E.**, Northeast, D. B., Jin, J., Laferrière, P., Korkusinski, M., Poole, P. J., Williams, R. L., & Dalacu, D. (2023). On-chip indistinguishable photons using III-V nanowire/SiN hybrid integration. *Phys. Rev. B*, *108*(19), 195417,
2. **Yeung, E.**, Sorensen, K., Northeast, D. B., Milanizadeh, M., Poole, P. J., Williams, R. L., Dalacu, D. (2026). Nanowire Quantum Dot Emitters Evanescently Coupled On-Chip via Curved and Segmented SiN Waveguides. *ACS Photonics*, *13*(15), 4185-4192,
3. **Yeung, E.**, Zeeshan, M., Northeast, D. B., Poole, P. J., Williams, R. L., Dalacu, D. (2026). Plug-and-play single-photon sources using chip-integrated nanowire quantum dots. *Appl. Phys. Lett.* *129*(8), 084005.

Detailed contributions are stated in the preamble for each paper presented in Ch. 4, Ch. 5, and Ch. 6.

In addition, I contributed equally to the following papers, summarized in Ch. 2, with co-author Patrick Laferrière to develop the nanowires as the required resources for the articles presented in this thesis:

1. Laferrière, P., **Yeung, E.**, Giner, L., Haffouz, S., Lapointe, J., Aers, G. C., Poole, P. J., Williams, R. L., & Dalacu, D. (2020). Multiplexed Single-Photon Source Based on Multiple Quantum Dots Embedded within a Single Nanowire. *Nano Letters*, *20*(5), 3688–3693,
2. Laferrière, P., **Yeung, E.**, Korkusinski, M., Poole, P. J., Williams, R. L., Dalacu, D., Manalo, J., Cygorek, M., Altintas, A., & Hawrylak, P. (2021). Systematic study of the emission spectra of nanowire quantum dots. *Appl. Phys. Lett.*, *118*(16), 161107,
3. Laferrière, P., **Yeung, E.**, Miron, I., Northeast, D. B., Haffouz, S., Lapointe, J., Korkusinski, M., Poole, P. J., Williams, R. L., & Dalacu, D. (2022). Unity yield of deterministically positioned quantum dot single photon sources. *Sci. Rep.*, *12*(1), 6376, and

4. Laferrière, P., Yin, A., **Yeung, E.**, Kusmic, L., Korkusinski, M., Rasekh, P., Northeast, D. B., Haffouz, S., Lapointe, J., Poole, P. J., Williams, R. L., & Dalacu, D. (2023). Approaching transform-limited photons from nanowire quantum dots using excitation above the band gap. *Phys. Rev. B.*, *107*(15), 155422.

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List of Publications

1. Laferrière, P., **Yeung, E.**, Giner, L., Haffouz, S., Lapointe, J., Aers, G. C., Poole, P. J., Williams, R. L., & Dalacu, D. (2020). Multiplexed Single-Photon Source Based on Multiple Quantum Dots Embedded within a Single Nanowire. *Nano Letters*, *20*(5), 3688–3693.
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3. Wong, V., **Yeung, E.**, Cole-André, M., & Fanchini, G. (2022). Phase modulated scanning near field optical luminescence imaging as a probe of exciton dark lifetime and diffusivity in persistently luminescent micro- and nanocrystals. *J. Mater. Chem. C*, *10*(36), 12975–12985.
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6. **Yeung, E.**, Northeast, D. B., Jin, J., Laferrière, P., Korkusinski, M., Poole, P. J., Williams, R. L., & Dalacu, D. (2023). On-chip indistinguishable photons using III-V nanowire/SiN hybrid integration. *Phys. Rev. B*, *108*(19), 195417.
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10. Kamel, N. G., Mansourzadeh, A., Gautam, U., Kavatamane, V. K., Singh, A., **Yeung, E.**, Northeast, D. B., Barclay, P., Poole, P. J., Dalacu, D., & Oblak, D. (2026). Erbium-Doped Fibre Quantum Memory for Chip-Integrated Quantum-Dot Single Photons at 980 nm. *npj Quantum Inf.*

11. Singh, P., Jagde, J. K., Jain, M., **Yeung, E.**, Northeast, D. B., Moisa, S., Mohammed, S. J., Lapointe, J., Rajnis, U., Kienast, A., Poole, P. J., Dalacu, D., & Hall, K. C. (2026). Quantum Emission in Monolayer WSe2 Transferred onto InP Nanowires. *Adv. Optical Mater.*, *14*(3), e02992.
12. **Yeung, E.**, Sorensen, K., Northeast, D. B., Milanizadeh, M., Poole, P. J., Williams, R. L., Dalacu, D. (2026). Nanowire Quantum Dot Emitters Evanescently Coupled On-Chip via Curved and Segmented SiN Waveguides. *ACS Photonics*, *13*(15), 4185-4192.
13. **Yeung, E.**, Zeeshan, M., Northeast, D. B., Poole, P. J., Williams, R. L., Dalacu, D. (2026). Plug-and-play single-photon sources using chip-integrated nanowire quantum dots. *Appl. Phys. Lett.* *129*(8), 084005.

*To my parents,
for all their love and support throughout all my years of school*

Thank you



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Nomenclature

T_1 lifetime

X neutral exciton

X^- negatively charged exciton

XX biexciton

$g^{(1)}(\tau)$ first-order correlation measurement

$g^{(2)}(\tau)$ second-order correlation function

T_2 coherence time

P_{sat} power at saturation

Acronyms

AES anisotropic exchange splitting. 17

AR anti-reflective. 26

BS beam-splitter. 40, 44, 45

CBE chemical beam epitaxy. 9, 30

CCD charge-coupled device. 36, 37

CMOS complementary metal-oxide-semiconductor. 33

CW continuous-wave. 15, 22, 26, 35, 44

DUT device under test. 28

FDTD finite-difference time-domain. xv, 8

FSR free spectral range. xvii, 37

FSS fine-structure splitting. 17, 37, 38

HOM Hong-Ou-Mandel. 44, 45

HWP half-wave plate. 26, 27, 44

InAsP indium arsenide phosphide. 6

InP indium phosphide. 6, 8

MZI Mach-Zehnder interferometer. 42, 44

NA numerical aperture. 31

NRC National Research Council of Canada. 14, 31

NW nanowire. 2, 7

NW-QD nanowire-quantum dot. 3, 6, 7, 12, 14, 16, 17, 19, 25, 30, 31, 90

PBS polarized beam-splitter. 26–28

PIC photonic integrated circuit. xvi, 24, 28, 32, 90, 91, 122, 125

PL photoluminescence. 10, 15, 30, 35, 37, 38

PM polarization maintaining. 26

QD quantum dot. 2, 4–7, 9–12, 14, 15, 17, 35, 44

QKD quantum key distribution. 1

SA-VLS selective-area vapour-liquid-solid. 8

SEM scanning electron microscope. xv, 8, 32

Si₃N₄ silicon nitride. 3

SiO₂ silicon oxide. 8

SK Stranski–Krastanov. 5, 6

SNSPD superconducting nanowire single-photon detector. 34, 37, 125

SPAD single-photon avalanche diode. 34, 37

TE transverse electric. 26

TEM transmission electron microscope. 9

TM transverse magnetic. 26

TPI two-photon interference. 44

TRPL time-resolved photoluminescence. 39, 42

VOA variable optical attenuator. 38

Chapter 1

Introduction

Single-photons are a necessary building block for many quantum information technologies [2] including quantum computing [3], quantum communications [4], and quantum cryptography [5]. In these applications, information is encoded in the quantum states of individual photons. Particularly, in quantum cryptography protocols such as quantum key distribution (QKD), the security of communication relies on transmitting information through individual photons that cannot be intercepted or cloned without detection [6]. As such, an ideal single-photon source needs to be able to generate single-photons in a controlled, reproducible, and highly pure manner to ensure reliable quantum operations. In other words, the single-photons need to be monochromatic, have a lifetime-limited linewidth, have a long coherence time, and be indistinguishable.

An ideal single-photon source should satisfy the following criteria [7, 8, 9]:

- **brightness:** have high collection of all photons produced; typically defined as the number of photons collected per excitation pulse into the first lens [10],
- **high-repetition rate:** high photon emission rates for fast quantum operations,
- **single-photon purity:** each emission event should produce exactly one photon, with minimal multi-photon probability (quantified by $g^{(2)}(\tau=0) = 0$),
- **spectral purity:** emitted photons should have stable phase and narrow spectral linewidths, and
- **indistinguishability:** each photon is identical in every way (i.e., polarization, spatial mode, and temporal profile) to allow for perfect quantum interference.

Single-photon generation should ideally be deterministic where emission occurs at controllable, user-defined times (i.e., “on-demand”). However, these requirements are meaningless without an efficient method to couple single photons into desired quantum channels [7] such as optical modes compatible with fibre transmission or on-chip waveguides. The source should also be stable and maintain its optical properties under varying environmental conditions. Ultimately, the source needs to be scalable for implementation into larger quantum networks [11].

1.1 Single-photon generation: semiconductor quantum dots

Single-photons can be generated using either probabilistic or deterministic methods. Probabilistic sources include nonlinear optical processes such as parametric down-conversion and four-wave mixing, where a high-energy photon is split into two lower-energy photons [12]. In these cases, the generation of single photons is random, and multi-photon events cannot be completely suppressed, hence limiting their efficiency. Deterministic sources, on the other hand, generate single photons on-demand. Examples include single atoms, ions, molecules, 2D materials, colour centres in crystals, and semiconductor quantum dots [8, 13]. These sources emit a photon each time an excitation pulse is applied, thus enabling controlled and repeatable operation.

In solid state systems, single photons are typically generated when an excitation laser promotes an electron to a higher energy level within a quantum dot [14, 13]. When the electron recombines with a hole, the energy difference is released as a single-photon. This process is known as photoluminescence. To make quantum information technologies compatible with our current telecommunication networks, we want to use a material system capable of generating single-photons reliably in the telecom O-band (1260 ~ 1360 nm) [15] and C-band (1530 ~ 1565 nm) range [16] such as semiconductor quantum dots.

Semiconductor quantum dots (QDs) are often referred to as “artificial atoms” because they confine charge carriers (electrons and holes) in all three dimensions, thus creating discrete energy levels similar to those in natural atoms. QDs are highly appealing as single-photon sources because their emission can be tailored from the infrared to telecom ranges by adjusting their material composition, size, and strain environment [8]. They can also be integrated directly into semiconductor fabrication processes [17] which enable scalable integration into compact devices [11]. The emission can also be electronically or optically controlled and tuned using temperature [15, 18], strain [19, 20], or external fields [21].

1.1.1 III-V nanowires embedded with quantum dots

QDs are formed from a strain driven process due to the lattice mismatch between the QD material and underlying layer. This results in the QDs being randomly nucleated on the substrate. The dots are subsequently embedded in the host semiconductor material which is typically a high III-V high-index semiconductor. Due to the total internal reflection, it is very difficult to get appreciable emission out of the material into free-space from this bulk format.

Embedding the QDs within nanowire (NW) waveguide structures will make the QD emit into the fundamental HE_{11} mode of the nanowire. In so doing, the nanowire geometry can be designed to optimize the single-photon collection efficiency to an objective and force the emission to propagate in one-dimension along the NW core. The NW structure also ensures the minimizing of losses from total internal reflection that plagues self-assembled Stranski–Krastanov QDs in a bulk host semiconductor [22, 23].

Unlike cavity-based systems that require spectral matching [24, 25], nanowires support broadband emission suitable for various quantum dot excitonic states [26]. Another key advantage of nanowire-quantum dots (NW-QDs) is that they can be grown in patterned arrays using selective-area vapour-liquid-solid (SA-VLS) epitaxy [1]. This bottom-up growth ensures the QD is positioned at a predetermined location along the length of the nanowire. It also allows for multiple QDs to be grown within the same nanowire with adequate vertical spacing [27]. By growing the nanowires in arrays of sufficiently large pitch, it is possible to excite and collect the emission from only one NW-QD at a time. The patterning also allows for quick and easy identification and relocation of a nanowire within a sample for a future study.

1.2 On-chip integration for scalability

The free-standing nanowires on the growth substrate are limited to free-space optical excitation and collection schemes. However, the in-plane generation and manipulation of single-photons is necessary for integrating into photonic integrated circuits. In order to achieve this, the nanowires need to be transferred from the growth substrate to a chip containing the photonic integrated circuitry.

The potential nanowire candidates for transfer are pre-characterized prior to integration into a quantum photonic circuit. The tapering of the nanowire allows for the evanescent coupling of the QD emission into an adjacent silicon nitride (Si_3N_4) waveguide [28].

This thesis focuses on developing a pick-and-place [29] platform to fabricate on-chip quantum integrated circuits whereby III-V quantum dot sources are integrated with silicon-based integrated circuits. The long-term goal of this research is to develop a stable and scalable source of single photons suitable for plug-and-play applications.

1.3 Outline of thesis

A basic background of the concepts needed to understand this thesis is presented in Chapter 2. The experimental methods used to investigate the concepts discussed in the previous chapter are presented in Chapter 3. Chapter 4 is a proof-of-concept demonstration that our hybrid integration approach is a good source for in-plane single-photon generation. In order to build and design advanced architectures, Chapter 5 explores advanced architectures for more efficient single-photon generation. Ultimately, the goal is to have a plug-and-play source of single-photons that can be placed into any cryostat and operated without active alignment. We achieve this in Chapter 6 by packaging the chips with fibre-bundles. The conclusion and future works are presented in Chapter 7.

Chapter 2

Background

In this chapter, a brief background on the quantum light sources studied in this thesis is provided, namely III-V nanowires embedded with quantum dots (NW-QDs). An overview of the optical properties of the NW-QDs is presented based on Refs. [30] and [31]. In each of these papers, I contributed equally to the data acquisition, analysis, and writing of the original draft with my coauthor P. Laferrière, who was the main lead in each project. These papers form the basis for developing the nanowires used in the quantum photonic circuits presented in Ch. 4 to 6 of this thesis. The latter half of the chapter provides a general overview of the basic concepts pertaining to photonic waveguides.

2.1 Quantum confinement in quantum dots

Quantum dots (QDs) can be understood as the physical realization of the “particle in a box” model. Within a QD, charge carriers such as electrons and holes are spatially confined in all three dimensions. Quasi-particles called excitons, which are bound electron-hole pairs, with a confinement scale on the order of the exciton Bohr radius. As the size of the quantum dot decreases, the energy levels available to the carriers become more widely spaced as illustrated in Fig. 2.1 (a). In the context of the particle-in-a-box model, these energy levels can be described by

$$E_n = \frac{n^2 \hbar^2 \pi^2}{2mL^2} = \frac{n^2 h^2}{8mL^2}, \quad (2.1)$$

where m is the mass of the particle in the box, L is the length of the box, n is an integer index of such a level, and the reduced Planck’s constant is $\hbar = h/2\pi$. From Eq. 2.1, it can be seen that as L decreases, the relative difference between the energy levels increases. An example of how the energy levels change with varying lengths is shown in Table 2.1 for $n = 1$ and $E_1^* = h^2/8m$.

From the table, it is clear that reducing the size of the box by half causes the energy levels to become four times higher. A similar phenomenon occurs in quantum dots. As their size decreases, the spacing between energy levels increases. These discrete and quantized

L	$n = 1$	$n = 2$	$n = 3$
0.5	$4E_1^*$	$16E_1^*$	$36E_1^*$
1	$E_1^* = \frac{\hbar^2}{8m}$	$4E_1^*$	$9E_1^*$
2	$\frac{1}{4}E_1^*$	E_1^*	$\frac{9}{4}E_1^*$

Table 2.1: Discrete energy levels for particle in a box.

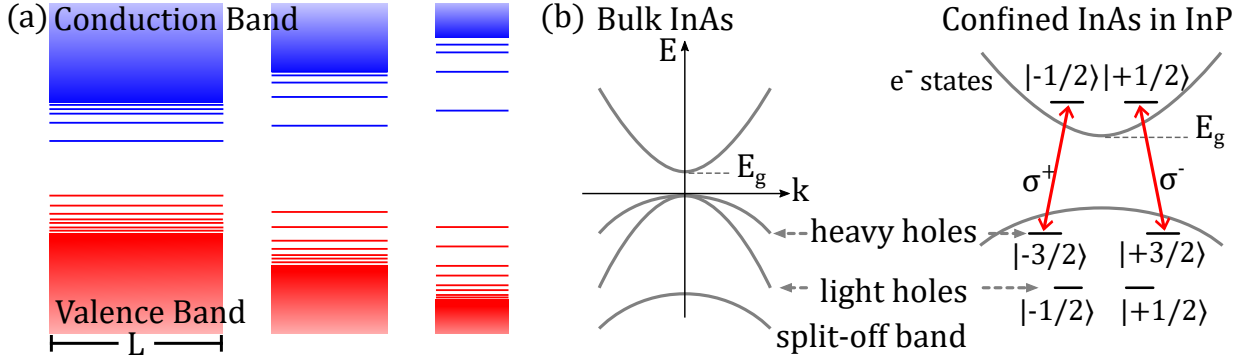


Figure 2.1: Confinement of a quantum dot. (a) Schematic of quantum confinement effects. The spacing between discrete energy levels increases with decreasing size. (b) E-k diagram of bulk InAs (left) and confined InAs in InP (right).

energy levels are similar to those in atoms, which is why quantum dots are often referred to as “artificial atoms” [9].

There are three main types of quantum dots including gate-defined quantum dots, colloidal quantum dots, and epitaxial quantum dots. Gate-defined quantum dots use electrical fields to create a confining potential. Colloidal quantum dots consists of nanoparticles suspended in a solution. Epitaxially grown quantum dots are grown directly on a substrate. For the last, confinement is achieved by embedding a small band gap semiconductor (e.g., InAs) within a larger band gap host material (e.g., InP). In particular, Stranski–Krastanov (SK) quantum dots are self-assembled due to the lattice mismatch between the substrate and the dot material. These QDs are usually a few nanometers tall, thus making them strongly confined in the growth (z) direction, and this confining potential is typically assumed to be parabolic. Confinement in the xy directions are less strict, and the confining potential is more abrupt. The energy levels within such a quantum dot can be expressed as [32]

$$E = (n_x + n_y + 1)\hbar\omega_0 + \left(\frac{\pi^2\hbar^2}{2m^*L^2}\right)n_z^2, \quad (2.2)$$

where n_x and n_y are integers ≥ 0 , n_z is an integer ≥ 1 , L is the height of the quantum dot, ω_0 is the energy-level spacing, and m^* is the effective mass of either the electron or the hole.

When these energy levels are populated with charge carriers and recombination occurs,

photons are emitted with energies described by [33]

$$E = E_e + E_h - V_{eh}, \quad (2.3)$$

where the energy of the electron, E_e , and energy of the hole, E_h , are defined by Eq. 2.2 using their respective effective masses, and V_{eh} represents the Coulomb interaction between the carriers, which corresponds to the exciton binding energy for a single electron-hole pair.

2.2 III-V nanowires embedded with quantum dots

Since the QD can emit in any direction, only a small percentage of the emission from an SK dot is collected [22]. To mitigate the poor collection efficiency, the QDs can be put in different photonic structures [8] such as cavities [34, 35] and waveguides [1, 24, 36]. These structures can be designed for single-mode operation for efficient coupling to optical fibres.

QDs in cavities (e.g., micropillars) are made by sandwiching the QD between two stacks of Bragg reflectors consisting of high and low index layers [37]. Since discrete modes exist within the cavity structure, these structures require careful spectral and spatial matching of the QD emission mode to the cavity modes for optimal emission. To optimize collection of the QD emission, the cavities can be tuned (e.g., with temperature or strain) and operated under resonant excitation [38]. Placing a QD in a cavity takes advantage of the Purcell effect to enhance the spontaneous emission rate and generate highly indistinguishable photons. Conversely, QDs in waveguides are based on propagating modes in a continuum. The spectral matching tolerances of QDs in waveguides are relaxed, thus allowing for a broader spectral range of supported wavelengths. Above-band excitation can be used to optically excite the QD and collect the full emission spectrum from the QD. Although emitted photons will not have as high a degree of indistinguishability compared to QDs in cavity, the broad range of supported wavelengths within the waveguide is a key advantage that allows for the study of different emission complexes within the QD which is of particular interest in this thesis.

The NW-QDs used in this thesis contain indium arsenide phosphide (InAsP) quantum dots in bulk indium phosphide (InP) nanowires. Bulk InP has a band gap of 1.34 eV at 300 K (1.42 eV at 4 K) [39], and bulk indium arsenide (InAs) has a band gap of 0.354 eV at 300 K [40] (0.42 eV at 4 K [41]). Since InP has a much larger band gap than InAs, InAs makes a good candidate as a quantum dot in a bulk InP material.

The band gap of the quantum dot is determined by the bulk band gap of the dot material (see Fig. 2.1 (b, left)) which will depend on the composition of the dot and the lattice mismatch-mediated strain between the dot material and the host material (also dependent on composition). We note that for strong confinement in the growth direction, the degeneracy of the heavy-light hole bands is removed (see Fig. 2.1 (b, right)) with the heavy hole band forming the ground state [9]. The heavy hole and light hole states have a split-orbit energy difference of 0.41 eV [40] for an InAs QD. As such, this thesis will limit the focus to just the heavy hole band. The description above provides a simplified view of

the electronic structure of a quantum dot. For a more detailed and realistic model, refer to Ref. [42]. There will be a more thorough discussion of the states in Ch. 2.2.4.

There are two main nano-antenna sources: top-down and bottom-up sources. Top-down nano-antennas (often referred to as micropillars [43, 24]) are obtained by etching a bulk material containing the QD. The number of emitters per micropillar can be controlled to a limited extent by adjusting the emitter density in the host material before etching [26]. Bottom-up sources (often referred to as nanowires) have both the emitter and waveguide defined through growth [1]. The advantage of a bottom-up technique is the ability to define exactly the desired number of QDs within a single nanowire with unity yield [31].

Nanowire (NW) waveguides are designed to optimize the light collection efficiency from an embedded quantum dot by coupling the emitter to the easy-to-collect fundamental HE_{11} optical modes compatible with collection through an objective and optical fibre [44, 45]. The higher refractive index of the nanowire waveguide ($n_{\text{InP}} \approx 3.37$ at $\lambda \approx 950 \text{ nm}$ [46]) compared to its surroundings allows light travelling through the nanowire to go through total internal reflection at the waveguide edges.

Fig. 2.2 (a) shows a typical tapered nanowire embedded as a thin segment of InAs along the core of the NW (see TEM image in inset of Fig. 2.3). After the quantum dot emission couples into the fundamental HE_{11} modes of the nanowire core, the mode propagates without spreading as it traverses the length of the nanowire waveguide. To limit the divergence of the beam as the light leaves the nanowire and continues into the far-field, the nanowire is typically tapered by $1\text{-}2^\circ$ [47] towards the tip. This expansion of the mode is shown as an overlay in Fig. 2.2. The taper also helps reduce possible back-reflection from the tip of the nanowire [48]. A top-down view of the HE_{11} mode profile is shown in the inset of Fig. 2.2 (b).

The expansion of the HE_{11} mode is also useful for evanescent coupling between the nanowire mode and another waveguide [11, 18]. A more detailed discussion will be presented in Ch. 4.

The efficiency of the source (η), is the fraction of photons emitted from the dot per excitation pulse collected at first lens by an external optical system. Assuming the dot emits one photon for each excitation pulse, the source efficiency will thus depend on the the fraction of spontaneously emitted photons from the QD into the nanowire fundamental HE_{11} modes with efficiency $\beta_{\text{HE}_{11}}$. The source efficiency is thus give by [1]

$$\eta = \alpha \beta_{\text{HE}_{11}} \eta_{\text{IQE}}, \quad (2.4)$$

where α is the fraction of photons coupled to the external optical system and η_{IQE} is the internal quantum efficiency of the emitter. For our NW-QDs system, α is 50% since half the emitted photons are lost due to emission in the opposite direction. One way to recuperate the lost photons is to use a back mirror deposited on the bottom of the nanowire [26, 49, 50]. Another way to recuperate the lost photons is to couple a nanowire embedded with a quantum dot located midway along its length to a curved waveguide. This approach will be discussed in more detail in Ch. 5.

The spontaneous emission from the QD will couple to either the confined ($\Gamma_{\text{HE}_{11}}$) or leaky (γ) modes. $\beta_{\text{HE}_{11}}$ is thus given by the fraction of the spontaneous emission into the

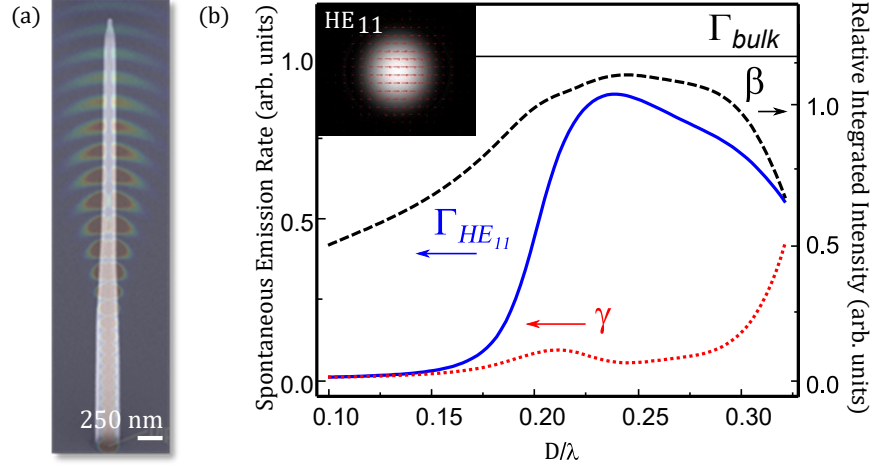


Figure 2.2: (a) Scanning electron microscope (SEM) image of a nanowire overlaid with an artistic representation of the expansion of the HE₁₁ mode as the light propagates through the taper is shown. (b) Finite-difference time-domain (FDTD) simulation showing the normalized spontaneous emission rates calculated for an InP waveguide at an operating wavelength of $\lambda = 950$ nm. Insert shows a simulation of the HE₁₁ mode in a nanowire waveguide in a top-down view. Lines are vectors indicating the direction of the electric field of one of the two degenerate polarization states. Adapted from [1].

fundamental mode relative relative to the emission in all modes represented as

$$\beta_{HE_{11}} = \frac{\Gamma_{HE_{11}}}{\Gamma_{HE_{11}} + \gamma}. \quad (2.5)$$

A finite-difference time-domain (FDTD) simulation of the emission rates $\Gamma_{HE_{11}}$, γ , and $\beta_{HE_{11}}$ as a function of the normalized waveguide diameter (D/λ), is shown in Fig. 2.2 (b) [1]. From the simulations, a peak efficiency of $\beta_{HE_{11}} = 95\%$ can be obtained for $D/\lambda \sim 0.25$. These calculations assume that the QD is centred laterally in the nanowire since any off-positioning will result in a reduction of the coupling into the fundamental HE₁₁ modes and potential coupling into higher order modes.

2.2.1 Growth

The nanowires are fabricated with a site-controlled, bottom-up approach using selective-area vapour-liquid-solid (SA-VLS) epitaxy. This approach is used to ensure that the desired number of quantum dots is grown within a nanowire core [51, 52, 27, 49, 53]. The schematic for the nanowire growth is presented in Fig. 2.3.

Starting with an indium phosphide (InP) wafer, the substrate is first coated with a 20 nm layer of silicon oxide (SiO₂) then a layer of resist (ZEP-A N-Amyl Acetate). Electron-beam lithography is used to pattern the ZEP resist which is transferred to the SiO₂ layer on the InP substrate via dry etching. The diameter of the SiO₂ mask openings define the final diameter of the nanowire waveguide.

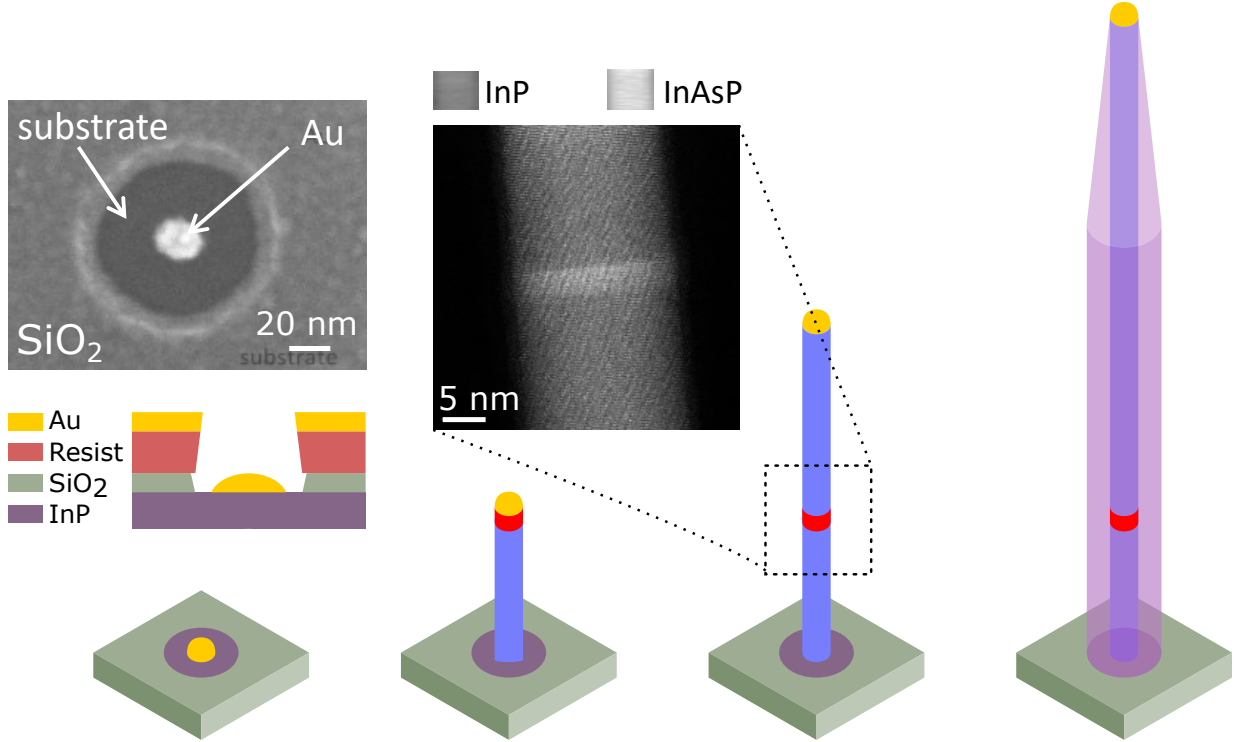


Figure 2.3: Cartoon showing the nanowire growth process. First column shows an SEM image of an InP wafer etched to expose a circular area of the InP substrate. Gold is deposited to the centre. The stacks of each layer prior to the lift-off process for removal of the resist and top Au layer are shown. Middle panel shows a TEM image of a QD segment consisting of ~ 10 monolayers of InAsP within a nanowire.

Wet etching with buffered oxide etch (BOE) containing hydrofluoric acid (HF) is used to expose the InP surface within the oxide holes. Fifty angstroms of gold (Au) is then deposited through electron beam evaporation. This thickness is chosen to achieve Au nanoparticles of 20 nm diameter inside the mask openings which will dictate the diameter of the InP core and InAs quantum dot. These steps correspond to the first panel of Fig. 2.3. A lift-off process removes the resist layer along with the unwanted Au layer while the Au particle inside the holes are intact.

The nanowires are grown with chemical beam epitaxy (CBE). A low flux of trimethylindium (TMI) and pre-cracked phosphine (PH_3) is injected into the chamber at $420\sim 435^\circ\text{C}$ [54]. Axial growth is promoted at the gold catalyst to ensure layer-by-layer growth only at the metal-semiconductor interface [1] and not the unmasked areas of a III-V substrate [55]. This first step in growing the nanowire core serves as a template for the crystalline structure of the InP. The nanowire core growth is shown in panels 2 and 3 of Fig. 2.3.

The growth of the core is temporarily interrupted at predefined times to grow the quantum dot. The flux of TMI is stopped, and arsine gas (AsH_3) is instead used to form a thin layer of InAsP before switching back to TMI to continue growth for the rest of the core. A transmission electron microscope (TEM) image of the QD layer in the nanowire core is shown between the second and third panels of Fig. 2.3. The geometry of the quantum

dot is determined by the diameter of the core (~ 20 nm) and the height (< 5 nm) of the quantum dot layer within the core [56, 1]. It should be stated again that the strongest confinement is in the vertical/axial growth direction.

The second growth step puts a cladding around the core that increases the nanowire diameter from 20 nm to approximately 200 nm to 250 nm. The larger diameter is designed to support the optical mode into which the quantum dot emits as shown in Fig. 2.2(b). Without the cladding, the emission will spread into the radiative modes of the surrounding vacuum. More importantly, the non-radiative recombination centres associated with the dot-surface interface will quench the emission. A higher PH_3 flux and growth temperature promotes both axial and radial growth (see fourth panel of Fig. 2.3). This cladding has the same crystal structure as the core [57]. The taper is a consequence of the radial and axial growth of the cladding, and a combination of the two determines the degree of the taper [58]. This can be degree of tapering can be controlled as will be shown in Ch. 4 and Ch. 5.

Further details related to the growth chip's layout will be presented in Ch. 3.2.

2.2.2 Excitation regimes for quantum dots in nanowires

There is strong quantum confinement in the QD due to the bandgap difference between the InP and InAs along with the small size of the QD. This leads to the formation of quantized levels in the conduction and valence bands such as the ones shown in Fig. 2.1 (a) and Fig. 2.4. When an electron and a hole are simultaneously trapped inside of the QD, they form a bound state (i.e., exciton) due to Coulomb interactions. When they radiatively recombine, a photon is emitted. This process is called photoluminescence (PL).

Once trapped, the electrons and holes are not able to escape back into the bulk material without an external source providing extra energy. The QD levels can be populated with carriers using either optical or electrical excitation techniques. For electrical excitation, the QD is typically embedded in the intrinsic region of the P-I-N diode. An electric current is used to generate charge carriers which are free to move throughout the device which can then relax into the QD [59]. Optical techniques typically use a laser to optically excite and promote an electron from the valence band into the conduction band. There are three common optical excitation techniques, namely above-band excitation, quasi-resonant excitation, or fully resonant excitation (see Fig. 2.4).

In above-band excitation, the QD is optically excited with a pump laser of energy equal to or greater than the band gap of the bulk wurtzite (WZ) InP nanowire (≥ 832 nm at 4 K) (see Fig. 2.4 (a)). This is the most widely used and easiest excitation scheme. Charge carriers are generated within the vicinity of the QD whereby an electron is promoted from the bulk InP valence band to the conduction band. An electron-hole pair is created in the process where both are free to diffuse throughout the material. This pair may recombine within the bulk InP nanowire or become trapped in lower energy states. The carriers will naturally seek the lowest possible energy configuration and become trapped in the potential well of the InAsP quantum dot, which lies at a lower energy level than the surrounding InP material (see Fig. 2.5 (a)). The trapped electron and hole will relax into the s-shell of

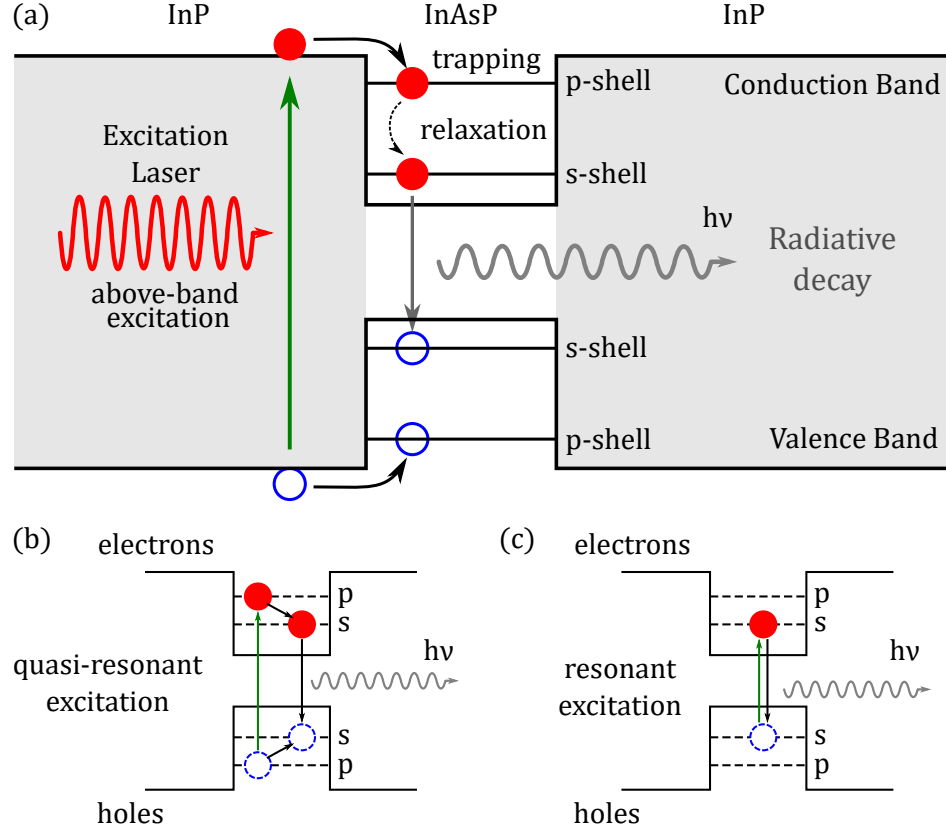


Figure 2.4: **Excitation regimes:** (a) above-band excitation, (b) quasi-resonant excitation, (c) resonant excitation. Electrons are shown as red circles, and holes are shown as empty blue circles.

the InAsP quantum dot and recombine to emit a photon with energy given by Eq. (2.3). This relaxation occurs over a few tens of picoseconds. The energy released as the carriers relax from higher energy levels to the s-shell is dissipated into the solid-state lattice as phonons (thermalization). As a result, this excitation scheme generates a lot of charge noise in the system [60] which inhomogeneously broadens the emission lines of the QD and causes decoherence of the emitted photons. A more thorough discussion on the effects of charge noise will be presented in later chapters.

For both quasi-resonant and resonant excitation, the QD is optically excited using a pump laser of an energy that matches one of the energy levels of the QD. Fig. 2.4 (b) shows an example schematic of quasi-resonant excitation whereby an electron is promoted from the p-shell of the valence band into the p-shell of the conduction band. Due to symmetry and the optical selection rules, only recombination of electrons and holes in the same shells is possible. For example, electrons in the p-shell orbital can only recombine with a hole in the p-shell, and an electron in the s-shell can only recombine with a hole in the s-shell [9]. Since the electrons and holes prefer to relax into the lowest energy state, the electron (hole) in the higher-order shell will relax into the s-shell (ground state) before it recombines with a hole also in the s-shell. During this relaxation to the ground state, the

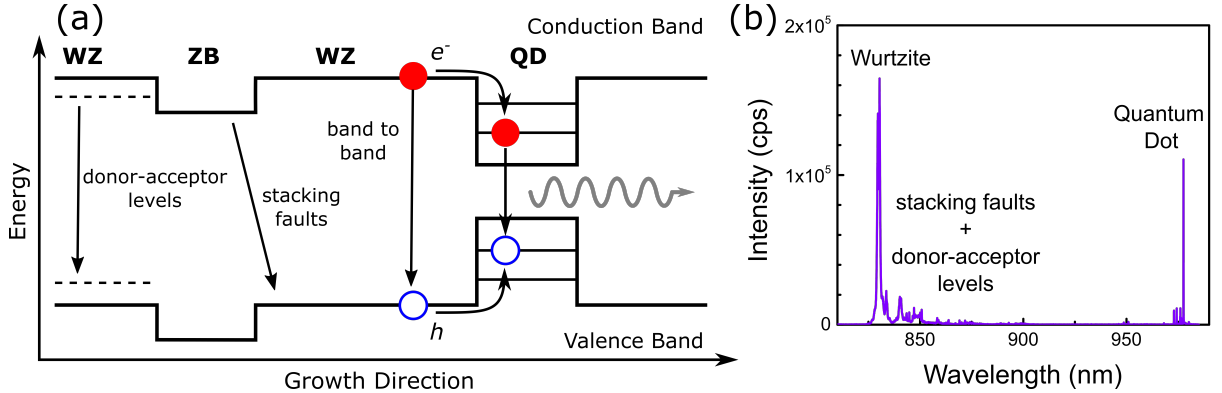


Figure 2.5: (a) Emission from an InP nanowire with an InAsP quantum dot including defects from stacking faults and donor-acceptor impurities. (b) A typical photoluminescence spectrum of a nominal nanowire-quantum dot.

excess energy is released as phonons. Since these methods directly pump into the energy levels of the QD, both quasi-resonant and resonant excitation schemes greatly reduce the thermalization caused from flooding the entire nanowire with carriers as is in above-band excitation. Thermalization still occurs as the p-shell relaxes into the s-shell, but this is relatively small compared to above-band excitation.

Strictly resonant excitation uses a laser tuned exactly to one of the emission energies of the QD in the s-orbital (i.e., exciton or charged exciton) [61]. Fig. 2.4 (c) shows the schematic for strictly resonant excitation. Unlike quasi-resonant excitation, there is no relaxation from higher orbitals to the s-shell. Resonant pumping schemes allow for coherent manipulation of excitonic states without creating additional charge carriers and multi-excitonic states [62], thus minimizing dephasing processes. A key problem in resonant excitation schemes is the difficulty in separating the pump laser from the emission energies. In free-space setups, pump laser rejection is achieved through spatial and polarization filtering [63] to a high degree using a quarter-wave plate with a confocal microscope [64, 65]. Filtering out of the pump laser resonant with the QD emission energy is not trivial for an on-chip setup and is not used in this thesis.

2.2.3 Optical properties: photoluminescence

The typical emission spectrum of a NW-QD using above-band excitation is shown in Fig. 2.5 (b). Apart from the QD emission, there are multiple emission lines associated with the bulk InP material, namely the wurtzite, zincblende, stacking faults, and donor-acceptor levels seen at wavelengths less than 890 nm. Fig. 2.5 (a) shows the different emission processes that can occur in the InP/InAsP nanowire system.

Through the SA-VLS growth process, two crystal phases, called wurtzite (WZ) and zincblende (ZB), arise in the InP core [1, 66]. ZB is the more common crystal phase in bulk InP. However, the WZ crystal phase is the favoured growth during VLS [67] for the InP nanowires. WZ is the predominant emission from the nanowire [68] emitting at

$\lambda \sim 832$ nm at 4 K. ZB is purposely grown to be less common within the nanowire and emits at $\lambda \sim 875$ nm at 4 K. Typically the ZB emission seen in the nanowire PL spectrum is from the substrate.

Stacking faults consists of layers of the ZB phase inserted between the dominant WZ phase which give rise to quantum confinement of carriers along the length of the nanowire. Their emission energies lie between the two phases (832 $\sim$ 875 nm). If the stacking faults are within the vicinity of the quantum dot, charge traps are created which lead to spectral diffusion and decoherence of the quantum dot emission [9, 69] thus resulting in a broader emission linewidth for the QD. Impurities or lattice defects (missing In or P atoms) in the InP give rise to donor-acceptor levels which are seen as a band of emissions centred around 860 nm [70]. The typical emission wavelengths are summarized in Table 2.2.

Centre Wavelength (nm)	Source
832	Wurtzite band gap
840	Stacking faults in nanowire
860	Donor-acceptor band
875	Zincblende band gap
890-980	Quantum dot emission

Table 2.2: Commons sources of peaks on a PL spectrum from an wurtzite InP nanowire with embedded InAsP quantum dot grown on an InP substrate cooled to 4 K.

2.2.4 Excitons, charged excitons, and biexcitons

Within the s-shell, it is convenient to think of the spin states of the carriers in terms of their angular momentum $|J, J_z\rangle$. For III-V semiconductor bulk materials, the lowest level of the conduction band has a total angular momentum of $j = \pm 1/2$, whereas for the holes in the valence band, the total angular momentum can be either $j = 3/2$ or $j = 1/2$. Specifically, for the component projected onto the z -axis (growth direction), the $j = 3/2$ band consists of the heavy holes with $j_z = \pm 3/2$ and light holes with $j_z = \pm 1/2$. The light holes and heavy holes have a large separation between the them (see Fig. 2.1 (b) (right)). Since the complexes of interest are at the lowest energy levels, we will only be considering the heavy holes in the following discussion. The spin states for the electron and heavy holes can thus be written as [9]:

$$\begin{aligned}
 |e, \uparrow\rangle &= \left| \frac{1}{2}, +\frac{1}{2} \right\rangle; |e, \downarrow\rangle = \left| \frac{1}{2}, -\frac{1}{2} \right\rangle \\
 |h, \uparrow\rangle &= \left| \frac{3}{2}, +\frac{3}{2} \right\rangle; |h, \downarrow\rangle = \left| \frac{3}{2}, -\frac{3}{2} \right\rangle.
 \end{aligned} \tag{2.6}$$

There are therefore four possible spin states for the total angular momentum basis of the exciton state:

$$\begin{aligned} &|e \uparrow, h \downarrow\rangle \text{ and } |e \downarrow, h \uparrow\rangle \\ &|e \uparrow, h \uparrow\rangle \text{ and } |e \downarrow, h \downarrow\rangle \end{aligned} \quad (2.7)$$

In the case of the bright exciton, the spin-orientation of the charge carriers is anti-parallel ($|\uparrow, \downarrow\rangle$ and $|\downarrow, \uparrow\rangle$) which give a total angular momentum of $\Delta J = \pm 1$. These two excitonic states are typically energetically non-degenerate due to an exchange splitting resulting in a fine-structure splitting (FSS). This will be discussed in more detail later in Ch. 2.2.6. In contrast, dark excitons have parallel spin orientations of $|\uparrow, \uparrow\rangle$ and $|\downarrow, \downarrow\rangle$ and a total angular momentum of $\Delta J = \pm 2$. No recombination occurs in this latter case since it violates spin momentum conservation [8]. They are considered optically inactive until a spin flip occurs, allowing for recombination to occur [9]. There will be further discussion on the lifetimes of bright and dark excitons in Chapter 5.

Besides the neutral excitonic state, several additional many particle states are possible when extra electrons or holes are captured inside the QD. The possible “bright” variates of these states are shown in Fig. 2.6 (a). The neutral exciton state is denoted as $|X\rangle$ and consists of just the electron-hole pair. The biexciton state, $|XX\rangle$, consists of two electron-hole pairs and have a total angular momentum of $\Delta J = 0$. When one extra electron or hole is present, a charged exciton state (often referred to as a trion), $|X^{+/-}\rangle$, is formed. Recombination of a single electron-hole pair in any of these charge configurations produces a photon with a distinct emission energy. For the biexciton, the direct decay to the ground state is not possible due to optical selection rules. In the radiative cascade from the $|XX\rangle$ state to either of the two bright $|X\rangle$ states, a photon is emitted. Then another photon is emitted during the decay process from the $|X\rangle$ state to the QD ground state $|0\rangle$. This process is known as the XX-X cascade and is illustrated in Fig. 2.6 (b). The mixing of the $J_z = \pm 1$ states caused by the anisotropic exchange leads to two linearly polarized neutral-exciton states ($|X_H\rangle$ and $|X_V\rangle$). The photon generated by the recombination of the trion $|X^*\rangle$ state, on the other hand, is circularly polarized [24].

In this chapter, four different ways to identify each complex is presented: power-dependent photoluminescence, auto-correlation $g^{(2)}(\tau)$, cross-correlation $g^{(2)}(\tau)$, and high-resolution photoluminescence. The details on the experimental setup for each technique is presented in Ch. 3.

2.2.5 Determining excitonic complexes through photoluminescence spectra

In Laferrière, P., Yeung E., *et. al.*, *Appl. Phys. Lett.* (2021) [30], a systematic study on 42 different NW-QDs grown at the National Research Council of Canada (NRC) show that the majority of the NW-QDs showed three distinct emission lines corresponding to the X , XX , and X^- transitions. The emission spectra from four representative nanowires are shown in Fig. 2.7.

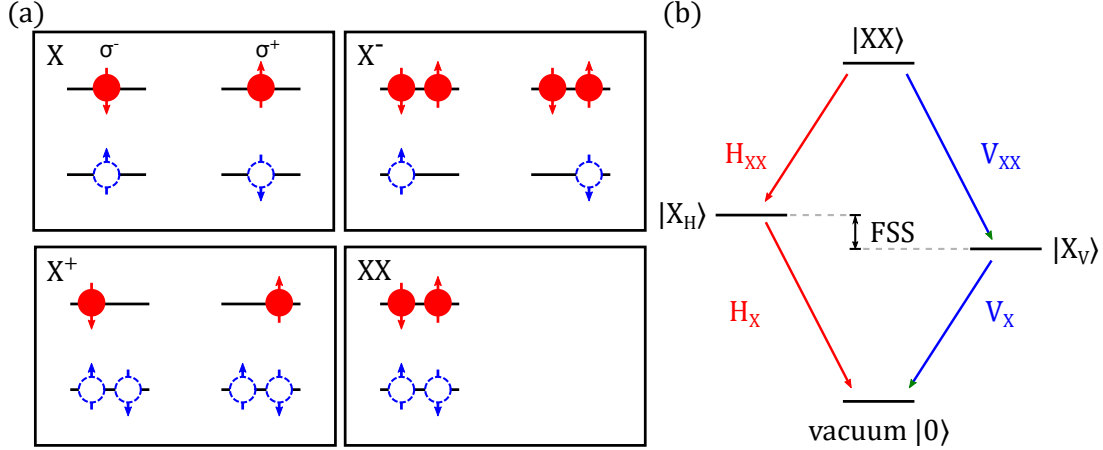


Figure 2.6: (a) Excitonic complexes of “bright” states. Top-left: neutral exciton, top-right: negatively charged exciton, bottom-left: positively charged exciton, bottom-right: biexciton. (b) Energy diagram for the biexciton-exciton (XX-X) cascade.

The nanowires are grown under the same conditions with the same nominal geometries (arsenic concentration of 18.2%, QD height of 4.1 nm, QD diameter of 18.2 nm), yet there are still differences in the emission spectra between the QDs (i.e., wavelength of complexes and spacing between complexes). This is due the random distribution of arsenic atoms in the dilute quantum dot [42].

To determine the identity of each complex, power-dependent photoluminescence (PL) spectra can be used to observe the rate at which the population of each complex increases relative to the excitation pump power at saturation of the X photons, P_{sat} .

Both the neutral exciton X and charged exciton X^- photons were observed at low pump powers. At much lower pump powers, the X photon should be the only one present [71]. This may not always be the case depending on the charge environment of the nanowire which may cause a shift in the Fermi level due to a slight doping of the InP which readily supplies the system with an extra charge carrier.

The XX emission count rates overtake both the X and X^- count rates as the pump power increases beyond P_{sat} . While the s-shell emissions were concentrated around 1.3 eV, hints of the p-shell (~ 1.35 eV) can be seen in the top row of the Fig. 2.7 when the system is pumped above-band at the P_{sat} of the X using a continuous-wave (CW) laser.

The lower panel of Figure 2.7 shows the count-rate power dependence for each of the three dominant peaks in a log-log plot. The exponent for each fit to the power-dependent peak intensity curve is denoted as n . Of the three emission lines, the XX transition is easily distinguished from the X and $X^{+/-}$ transitions as it will have the fastest rate of increase of the three, with $n_{XX} \sim 2$.

Based on a simplified Poissonian model [71], the neutral exciton is predicted to have $n_X \sim 1$ and the charged excitons to have $n_{X^*} \sim 1.5$ due to the presence of the extra charge carrier in the system. Experimentally, for the majority of the 42 QDs measured, it was possible to make a reasonable assumption on the identity of each emission line based on

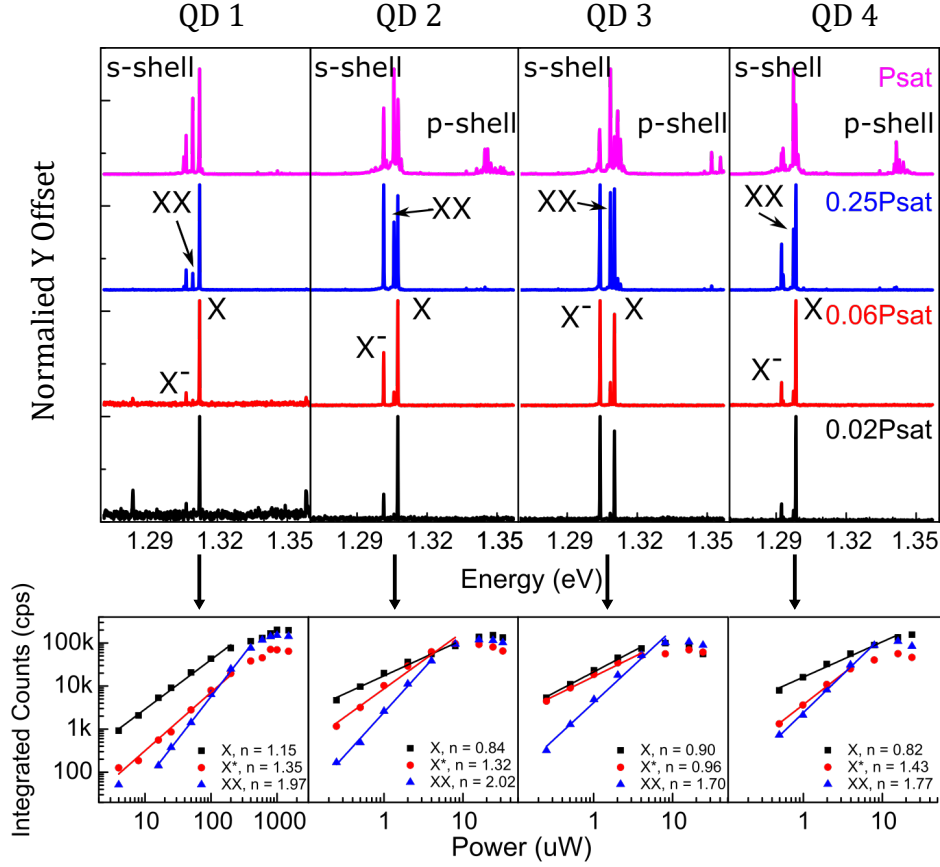


Figure 2.7: Power-dependent photoluminescence spectra of four quantum dots grown under the same conditions optically excited using an above-band pumping scheme. P_{sat} is the excitation pump power at saturation of the X transition.

this model.

A histogram of the extracted n values for each complex is shown in Fig. 2.8 (a). For the majority of the NW-QDs, a reasonable estimate could be made as to the identity of the three complexes, but a few NW-QDs had similar n values for the X and X^- transitions, such as the one in the third column of Figure 2.7 where $n_X = 0.90$ and $n_{X^-} = 1.07$. In such cases, measuring the cross-correlation between the XX and unknown complex was performed to correctly identify each transition. This will be discussed in a later section.

A spread of the emission energies from the X , XX , and X^- photons of the 42 different NW-QDs is shown in Fig. 2.8 (b). In this study, we found that the spread for the emission energy of X transition to be $E_X = 1301 \pm 5 \text{ meV}$, $E_{X^-} = 1296 \pm 5.6 \text{ meV}$ for the X^- transition, and $E_{XX} = 1300 \pm 4.3 \text{ meV}$ for the XX transition.

Computed transition energies using atomistic many-body calculations [30, 42, 72, 73] predicted the ordering in which the photons appear from highest to lowest energies to be X^+ , X , XX , and X^- . It should be noted that although the energies of the positively

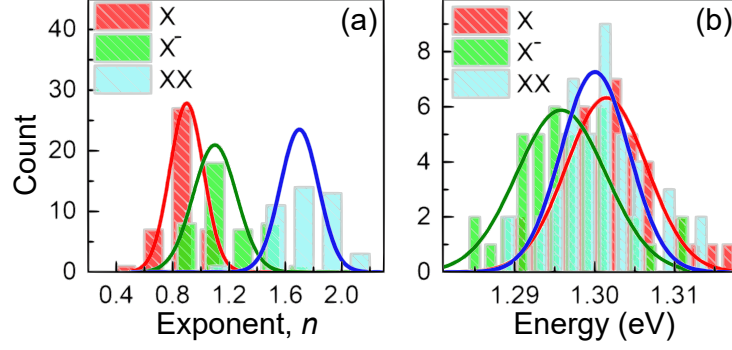


Figure 2.8: (a) Histogram of the exponential fit value, n . (b) Histogram for the energy distribution of different complexes extracted from 42 nanowire-quantum dots.

charged exciton, X^+ , were predicted by the model, they were not observed experimentally. In a later paper, Laferrière, P., Yeung, E., *et. al.*, *Sci. Rep.* (2022) [31], the PL of over a hundred NW-QDs were measured from a different growth sample. In this, a spread of 1264 ± 6 meV was measured for the X^- transition. Again, the X^+ was not observed. We suspect that there may be a charge trap in the vicinity of the QD supplying the QD system with a extra negative electron.

It should be noted that, here, we assumed the identity of the trion to be negatively charged based on its emission wavelength relative to the biexciton photons and neutral excitons photons as the model predicts. Identification of specific charge states requires an applied electric field to the structure [9, 74]. At the time, we did not have the capabilities to gate the free-standing nanowires.

The histogram in Fig. 2.8(a) shows that taking power-dependent spectra is a good starting point for identifying each transition, but a more reliable method is needed such as taking a high-resolution PL to observe the fine-structure splitting.

2.2.6 Fine structure splitting

Recall in Ch. 2.2.4, that the X and XX transitions have an anisotropic exchange splitting (AES)-induced fine-structure splitting (FSS) [63]. This splitting can be used to distinguish between X and X^- emission lines. Due to the high symmetry of the WZ nanowire quantum dot [42, 75], high resolution photoluminescence spectroscopy is needed to resolve this splitting. Both the X and XX transitions have anisotropic exchange splittings (AESs) whereas the charged excitons (X^*) does not [76].

Typical high-resolution PL spectra for the three transitions is shown in Fig. 2.9. It can be clearly seen that the X^* transition has only one peak, whereas the X and XX transitions have two peaks due to the FSS. The peak of each doublet corresponds to each orthogonally polarized emission (horizontal, H , and vertical, V) in the radiative decay process as shown in Fig. 2.6 (b). This is clearly seen through the use of a linear polarizer to filter out one of the two transitions. As the angle of the polarizer is rotated, the relative peak intensities exhibit a sinusoidal behaviour. No such sinusoidal behaviour is observed

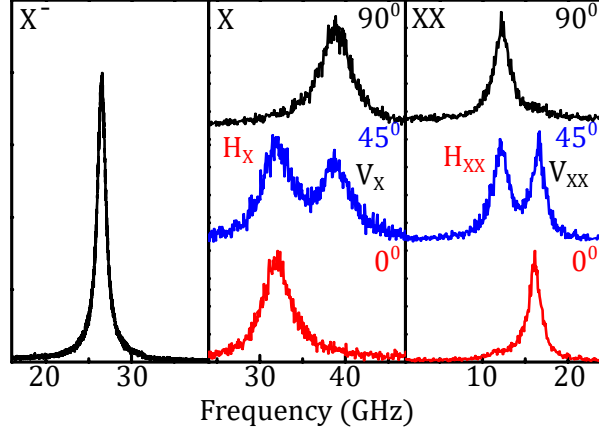


Figure 2.9: High-resolution spectra of the X^- , X , and XX of a nanowire-quantum dot using an etalon. For the X and XX transitions, the spectra are shown for three different orientations of a linear polarizer.

when tracking the peak emission of a charged exciton [76, 68, 77]. The FSS is typically in the few μeV range [78].

2.2.7 Radiative lifetimes and spectral purity

The radiative lifetime of a transition is how long it takes for a captured electron-hole pair to recombine and emit a photon. Each of the charge states recombine with a different probability [9]. A shorter radiative lifetime of an exciton is advantageous for maintaining the coherence of the emission by limiting the timescale at which spectral diffusion [69] and acoustic phonon-mediated pure dephasing [79] occur.

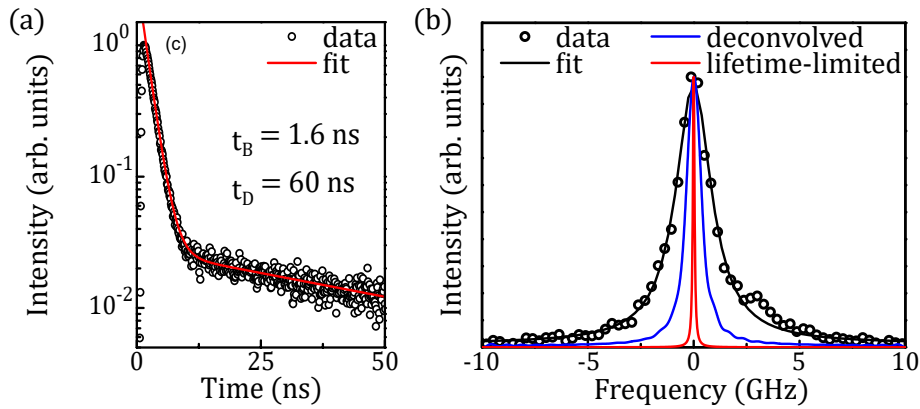


Figure 2.10: (a) Radiative lifetime of exciton complex. (b) Linewidth of the single exciton emission.

For the neutral and charged transitions, a double exponential decay function may be observed due to a faster bright state and a slower dark state. The slow component is due

to the forbidden dark transition consisting of parallel spins and a total angular momentum of $\Delta J = \pm 2$ (see Ch. 2.2.4). A spin-flip is required before recombination can occur to emit a photon. These two processes can be seen in Fig. 2.10 (a). For the charged exciton (see Fig. 2.6(a)), the dark state is due to a spin blockade as an electron or hole in the p-shell waits for a spin-flip before it can relax down to the s-shell. Conversely, the neutral biexciton exhibits a mono exponential decay in the lifetime due to the Pauli exclusion principle whereby the two electrons and two holes have antiparallel spins. As such, there is no wait period for a spin flip. More detail regarding the modelling will be presented in subsequent chapters.

For a two-level system, the radiative decay from the excited to ground state will have a spontaneous decay rate of $\Gamma_1 = 1/T_1$ where T_1 is the lifetime of the transition. Dephasing occurs as the phase relationships between quantum states are scrambled or lost at a rate of $\Gamma_\phi = 1/T_2^*$. The coherence of the single photon and impact due to dephasing mechanisms is expressed with [9]

$$\frac{1}{T_2} = \frac{1}{2T_1} + \frac{1}{T_2^*}, \quad (2.8)$$

where T_2 is the coherence time and T_2^* is the pure dephasing time. The factor of two arises before the T_1 term because the excited-state population, expressed as $P_e = |c_e|^2$, decays with a relationship of $|c_e|^2 \propto e^{-t/T_1}$, where c_e is the excited-state amplitude. As such, the population relaxation contributes a $\rho_{eg} \propto e^{-t/(2T_1)}$ where $\rho_{eg} = c_e c_g^*$ and c_g^* is the ground-state amplitude. In the ideal case where dephasing mechanisms are absent, $T_2/2T_1 = 1$. Fourier-transform limited photons have the narrowest possible linewidth determined by the emitter's lifetime, indicating that the photon may be highly coherent. There are two types of broadening that can affect the linewidth of the an emitter: homogeneous and inhomogeneous broadening. Homogeneous broadening is described by a Lorentzian distribution, and inhomogeneous broadening is described by a Gaussian distribution. The Fourier-limited lineshape is Lorentzian as shown in Fig. 2.10 (b) (red line).

Phonon-induced dephasing mechanisms are highly dependent on temperature, quantum dot size, and the structure embedding the QD. The confined carriers in the QD experience carrier-phonon interactions, carrier-carrier scattering, and nuclear spin noise [9]. Charge noise arise within the vicinity of the QD due to occupation fluctuations of the trapped charges affecting the local electric field [80, 81]. When the charge and spin fluctuations occur on a timescale greater than the radiative lifetime, it can lead to spectral wandering of the emission wavelength. These decoherent processes lower the indistinguishability of the emission.

In Laferrière *et. al.*, *Phys. Rev. B* (2023) [82], the linewidth of the X^- photon from several NW-QDs were measured and compared to the transform-limited linewidth. Typically, above-band gap excitation introduces charge noise and phonon interactions that broaden the emission line, reducing photon indistinguishability. A homogeneous linewidth of ~ 322 MHz for the X^- transition was measured for a NW-QD, which was around double the transform-limit predicted from the measured exciton lifetime of ~ 1 ns. It had a Lorentzian lineshape, indicative of homogeneous broadening. The linewidth remained stable at low excitation powers but broadened non-linearly as power increased beyond a

threshold ($P \geq 0.4P_{\text{sat}}$). Similarly, the linewidth showed little temperature dependence below 8 K, but increased rapidly above that temperature. Phonon-induced dephasing is therefore not the main source of residual broadening under low-power, low-temperature conditions. In that study, it was concluded that the careful engineering of the nanowire growth process and crystal structure can produce nearly transform-limited photons using above-band excitation.

2.2.8 Second-order correlation function

A second-order intensity correlation function, denoted as $g^{(2)}(\tau)$, is used to describe the degree of correlation or anti-correlation between photon detection events. More specifically, it describes the correlation between two temporally separated intensity signals with a time difference of $\tau = t_2 - t_1$. The classical second-order intensity correlation function is given by [83]

$$g^{(2)}(\tau) = \frac{\langle I(t)I(t+\tau) \rangle}{\langle I(\tau) \rangle^2} = \frac{\langle E^*(t)E^*(t+\tau)E(t+\tau)E(t) \rangle}{\langle E^*(t)E(t) \rangle^2}, \quad (2.9)$$

where $I(t) = E(t)E^*(t)$ is the intensity of the light at a given time t , and $I(t+\tau)$ denotes the intensity at some time t with a temporal delay of τ . The $\langle \dots \rangle$ denotes a time average.

In quantum optics, the electric field of a mode k can be rewritten in terms of photon creation, $\hat{a}^\dagger$, and annihilation, $\hat{a}$, operators using the relation [84]

$$\hat{E}_k(t) = \hat{E}_k^{(+)} + \hat{E}_k^{(-)} \quad (2.10)$$

with

$$\begin{aligned} \hat{E}_k^{(+)}(t) &\propto \hat{a}_k \cdot \exp(-i(\omega_k t - \vec{k} \cdot \vec{r})) \\ \hat{E}_k^{(-)}(t) &\propto (\hat{a}_k)^\dagger \cdot \exp(+i(\omega_k t - \vec{k} \cdot \vec{r})). \end{aligned} \quad (2.11)$$

For a single mode, the $g^{(2)}(\tau)$ function using the commutation relation can then be rewritten as

$$g^{(2)}(\tau) = \frac{\langle \hat{E}_k^{(-)}(t)\hat{E}_k^{(-)}(t+\tau)\hat{E}_k^{(+)}(t+\tau)\hat{E}_k^{(+)}(t) \rangle}{\langle \hat{E}_k^{(-)}(t)\hat{E}_k^{(+)}(t) \rangle^2} \quad (2.12)$$

$$\stackrel{(\tau \rightarrow 0)}{=} \frac{\langle (\hat{a}_k)^\dagger (\hat{a}_k)^\dagger \hat{a}_k \hat{a}_k \rangle}{\langle (\hat{a}_k)^\dagger \hat{a}_k \rangle^2} = \frac{\langle n(n-1) \rangle}{\langle n \rangle^2}. \quad (2.13)$$

in which n denotes a photon number. In particular, at $g^{(2)}(\tau=0)$, we get the conditional probability of how likely it is to detect a second photon at the same time the first photon was detected. In other words, it tells us the temporal photon coincidences. Eq. 2.13 can be further simplified using the second factorial

$$\begin{aligned} \langle n(n-1) \rangle &= \sum_{n=0}^{\infty} n(n-1) \cdot P(n) \\ &= \sum_{n=0}^{\infty} n^2 \cdot P(n) - \sum_{n=0}^{\infty} n \cdot P(n) \\ &= \langle n^2 \rangle - \langle n \rangle, \end{aligned} \quad (2.14)$$

where $P(n)$ is the probability of observing n photons. The variance is defined as the weighted sum of the squared deviations of the photon number n with respect to the mean occupation $\langle n \rangle$ and is given by

$$\begin{aligned} (\Delta n)^2 &= \sum_{n=0}^{\infty} (n - \langle n \rangle)^2 \cdot P(n) \\ &= \langle n^2 \rangle - \langle n \rangle^2. \end{aligned} \quad (2.15)$$

Using Eq. (2.14) and Eq. (2.15), Eq. (2.13) is simplified to

$$\begin{aligned} g^{(2)}(\tau = 0) &= \frac{\langle n^2 \rangle - \langle n \rangle}{\langle n \rangle^2} \\ &= \frac{(\Delta n)^2 + \langle n \rangle^2 - \langle n \rangle}{\langle n \rangle^2} \\ &= 1 + \frac{(\Delta n)^2 - \langle n \rangle}{\langle n \rangle^2}. \end{aligned} \quad (2.16)$$

This gives the statistical properties of the radiation [85]. Thermal light has a higher probability of emitting more than one photon at a time, with the photon number probability distribution exhibiting a Bose-Einstein distribution. The variance for thermal light is given as $(\Delta n^2) = \langle n^2 \rangle + \langle n \rangle$ [9]. Plugging this into Eq. (2.16), it can be seen that for thermal light, $g^{(2)}(\tau = 0) = 2$. One can think about it in the extreme case where both the photon number and variance are large, so the second term in Eq. (2.16) tends towards 1, so $g^{(2)}(\tau = 0) = 2$. In other words, the photons are simultaneously registered at both detectors (known as photon bunching).

Coherent light, such as from a laser, exhibits a Poissonian distribution, where the variance is $(\Delta n^2) = \langle n \rangle$. From Eq. (2.16), it can be seen that for coherent light, $g^{(2)}(\tau = 0) = 1$. In other words, the light is completely uncorrelated.

Fock states exhibit a sub-Poissonian distribution, where the variance is $(\Delta n^2) < \langle n \rangle$, so $g^{(2)}(\tau = 0) < 1$, leading to a reduced probability of emitting two photons at the same time (known as photon anti-bunching). Specifically, for an ideal single-photon source, the mean number of photons is $\langle n \rangle = 1$, and the statistical distribution is a delta function, so the variance is $\Delta n = 0$. Plugging this into Eq. (2.16), $g^{(2)}(\tau = 0) = 0$ for an ideal single-photon source. This is why the shorthand for the single-photon purity is typically described as $1 - g^{(2)}(\tau = 0)$.

Single-photon purity

Pure single-photon emission refers to a high probability of having exactly one photon in a time slot. The ideal source should emit exactly one photon at a time into a given spatiotemporal mode on demand. This value can be quantified with using the autocorrelation function, $g^{(2)}(\tau = 0) \sim 0$. If $g^{(2)}(\tau = 0) > 0$, then more than one photon may be present. The experimental details related to how to measure a second order correlation will be presented

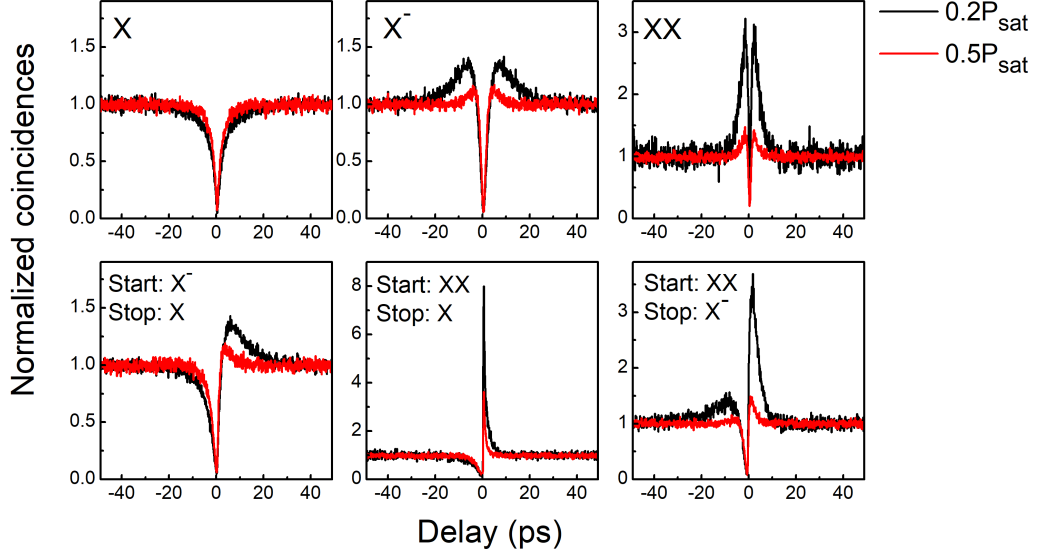


Figure 2.11: Autocorrelation (top) and cross-correlation (bottom) $g^{(2)}(\tau)$ measurements of X , X^- , and XX , respectively. Adapted from Laferrière, *et. al.*, *Appl. Phys. Lett.* (2021), **118**, 161107.

in Ch. 3.4.4. In this section, the differences in the $g^{(2)}(\tau)$ measurement for the X , X^- , and XX transitions are presented.

The top panel of Fig. 2.11 shows post-selected autocorrelation $g^{(2)}(\tau)$ measurements for the X , X^- , and XX photons, respectively, using above-band excitation with a CW laser. Each excitonic complex has its own signature autocorrelation and cross-correlation shapes [86, 87, 30]. Across all three complexes, a dip down to $g^{(2)}(\tau=0) \sim 0.05$ was observed, indicative of their single-photon purity.

In a two-level system, the excitonic state can be modelled by [87]

$$g^{(2)}(\tau) = 1 - e^{-(R+\Gamma)\|\tau\|}, \quad (2.17)$$

where R is the pump rate, and $\Gamma_1 = 1/T_1$ is the spontaneous decay rate. This model predicts that as the pump rate approaches zero, the half-width of the dip is the inverse of the radiative decay rate, $1/\Gamma$. This is seen for the X transition in Fig. 2.11 where as the pump power is lowered to $0.2P_{\text{sat}}$, the curve of the sidewalls widen.

For the X^- transition, the sidewalls of the curve do not change as a function of power. Instead, the X^- photons show bunching at short time delays. This bunching is associated with changes in the charge configuration of the dot, known as “blinking” [88, 89, 90]. It was mentioned in Ch. 2.2.5 that the presence of neighbouring traps within the vicinity of the quantum dot is supplying the QD with an extra electron. These traps may keep their populations substantially longer than the carrier reservoirs leading to secondary recombination [9]. As the pump power increases (i.e., the system is more readily filled with charge carriers), the time-scale associated with the blinking decreases. The lower the pump power,

the more pronounced the bunching. This bunching at lower powers is observed for the X^- transitions in other QD systems [91]. The degree of bunching is different from dot to dot.

The XX photon also shows strong bunching at short time delays as the power decreases. Strong “blinking” is typically observed in XX transitions [92].

Cross-correlation measurements for distinguishing excitonic complexes

The $g^{(2)}(\tau)$ is essentially the joint probability of detecting a photon at time $t = 0$ and another photon at time $t = \tau$. Cross-correlation measurements between two different photons, for example from the X and XX photons, can be used to help determine how likely a photon from one complex (e.g., X) is detected after a photon from a different complex (e.g., XX) is detected. This in turn tells how strongly correlated the two complexes are.

When the power-dependent intensities of the X and X^- photons yield similar exponential values, n , cross-correlation measurements can aid in the identification between the X and X^- photons. The lower panel of Fig. 2.11 shows the cross-correlation of the complexes from the top panel. The results from these correlation measurements are asymmetric due to the differing probabilities in detecting an event.

The lower first panel shows the cross-correlation between the X - X^- transitions. At short positive delays, there is a higher probability of detecting a X photon after the emission of a X^- photon. The dot is left with a single electron. A single hole is required for the QD to be ready to emit a X as opposed to another electron-hole pair needed for the X^- transition. This results in the bunching effect observed. This effect is washed out at higher pump powers because the system is more readily supplied with charge carriers. At short negative delays, after emitting a X photon, the dot is left empty, and hence no bunching is observed.

The lower middle panel shows a cross-correlation between the XX - X transitions. There is a strong correlation ($g^{(2)}(\tau=0) \sim 8$) between the two photons indicative of the $XX - X$ cascade [93, 94]. The XX photon are emitted first, followed by the X photon such that at $\tau = 0^+$, the emission is bunched, whilst at $\tau = 0^-$, the emission is anti-bunched. Once a XX photon is emitted, the system is ready to emit another X photon with a high probability. For close to $\tau < 0$, after a X photon is emitted, the system is emptied of charge carriers, and two more pairs of electron-holes are needed before another XX photon can be emitted. Similar observations were reported in [87, 93, 86].

Lastly, the lower right panel shows the cross-correlation between the X^- - XX transitions. At short positive delays, after the emission of a XX photon, the dot contains one electron-hole pair and requires the capture of an additional electron before a X^- photon can be emitted. Since we suspect there are traps near the dot readily supplying the QD with an electron, this is seen with a rather high probability. Note that this correlation is not as strong as with the XX - X cross-correlation because the QD still needs another electron before it is ready to emit a X^- photon. At short negative time delays, after the emission of a X^- photon, the dot is left with one electron and requires another electron and two holes to populate the dot before it can emit a XX photon.

2.3 Photonic Integrated Circuits

Optical waves travel in distinct modes through optical waveguides. In the case of a rectangular waveguide, a wave propagating along its length is confined by the difference in the refractive index between the waveguide material (Si_3N_4) and the surrounding material (SiO_2 or air). Waveguides are able to support wide bandwidths where the number of modes supported depends on the effective index, the dimensions of the waveguide, and the wavelength of the injected light. Rectangular waveguides will only propagate signals above a certain frequency, called the cut-off frequency. Below the cut-off frequency, waveguides fail to propagate light and the signal will be attenuated.

A well-engineered waveguide is essential for integrating pure single-photon sources and enabling the broader adoption of advanced photonic integrated circuits. Traditional free-space optical systems require frequent alignment, limiting their practicality outside laboratory environments. Transitioning to an all-fibre configuration enables efficient light transmission over extended distances. Compared to free-space setups, fibre-based photonic platforms offer improved stability, scalability, and portability. Further details on the design of an optical waveguide is presented in Ch. 6.

The photonic integrated circuits (PICs) used in this thesis were fabricated by AIM Photonics (Ch. 4), LIGEN-TEC (Ch. 5), and LioniX International (Ch. 6). Each silicon-based chip consisted of low pressure chemical vapour deposition (LPCVD) Si_3N_4 waveguides clad in SiO_2 .

Chapter 3

Experimental Methods

The methods used for characterizing photonic integrated circuits are presented in Ch. 3.1 which cover the setup and calibration of a fibre-based system, the measurement procedures, and the procedure to quantify of the propagation loss in a batch of chips with waveguides of identical dimensions. The layout of the nanowire samples is detailed in Ch. 3.2. Ch. 3.3 explains the assembly of hybrid III-V nanowire quantum photonic integrated devices using the pick-and-place technique. Finally, Ch. 3.4 provides an overview of the measurements used to characterize the optical properties of the quantum dots, including the emission spectrum, lifetime, single-photon purity, and indistinguishability of the single-photon source.

3.1 On-chip measurements

All measurements employed hybrid NW-QDs integrated onto Si_3N_4 photonic waveguide chips as the single-photon source. This section first describes the basic setup used to characterize the chips, including the system calibration process, followed by the procedure to measure the propagation loss.

3.1.1 Setup and calibration

Typically, two setups are used for chip measurements: one operating at room temperature and another located inside a cryostat. The room temperature setup is used to test the performance of different on-chip Si_3N_4 waveguide structures. Conversely, the cryostat setup is needed for measuring devices containing QDs which operate best below 4 K. Before conducting a series of transmission measurements on a chip, it is necessary to estimate the system loss.

For transmission measurements performed at room temperature, a chip is placed on a pedestal with a y-stage. Two lensed fibres on either side of the chip are used to couple light to and from the waveguides. The fibres are placed on two six-axis stages (x , y , z ,

roll θ_x , pitch θ_y , yaw θ_z) for precise alignment to the waveguide facets (see Fig. 3.1 (a) for schematic). Each fibre was a polarization maintaining (PM) 780HP spherical lensed fibre with an anti-reflective (AR) coating. The fibres are manufactured by Orbray to have a mode field diameter of $5.3 \pm 1.0 \mu\text{m}$ at 850 nm.

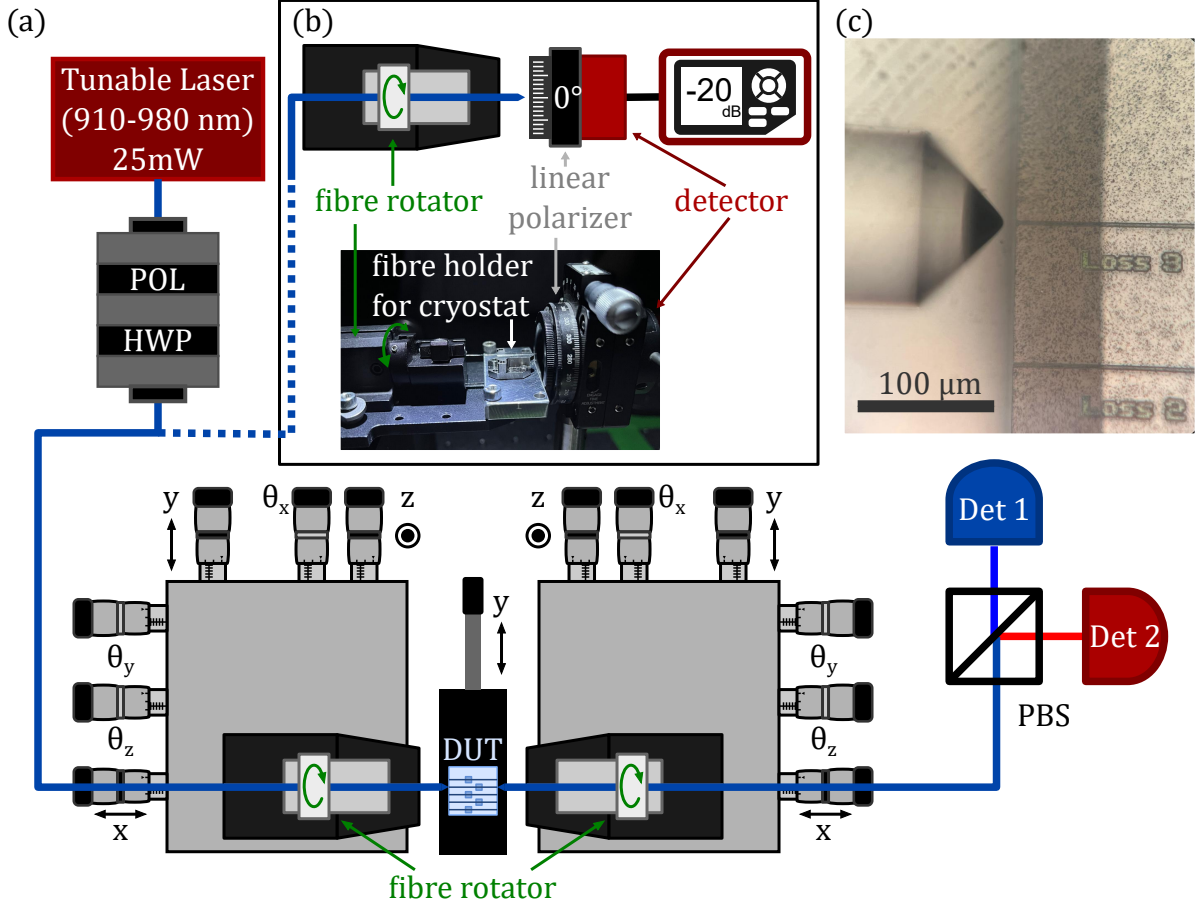


Figure 3.1: (a) Schematic for a typical fibre setup at room temperature using two 6-axis stages for the two fibres and a Y-axis stage to move the chip on a pedestal. (b) Schematic for the lensed fibre alignment. This procedure is done individually for both fibres. The photo shows the setup for aligning the lensed fibre on a cryostat compatible fibre holder. (c) Top-down optical image of a lensed fibre coupled to a waveguide.

A tunable CW laser from Toptica, capable of scanning across a 910 nm to 980 nm range, is used for the transmission measurements, which allows us to find the loss as a function of the wavelength. The output from the laser is linearly polarized and launched into the slow axis of a PM fibre. The ridge waveguides are designed to support the fundamental transverse electric (TE) and transverse magnetic (TM) modes for a band of wavelengths. As such, it is important to be able to launch light from the fibre such that the polarization is aligned to one of the two fundamental modes. In order to control the polarization through the fibre, an in-line half-wave plate (HWP) is used.

On the output side, the light collected from the second fibre is sent to a polarized beam-splitter (PBS), and its outputs directed to two fibre-coupled silicon photodiode de-

tectors (Thorlabs S151C). We can think of the two outputs as separating the input beam into the TE and TM modes. The connected power meters (Thorlabs PM103A) use a set measurement range (e.g., 3 mW) to deliver a proportional output voltage between -0.25 V to +2.5 V. The power meters are directly connected to the Toptica laser so that it can directly map the laser wavelength to the detected voltage. The voltage is converted back to a power reading using [95]

$$P = \frac{V_{\text{out}}}{2.5 \text{ V}} \times P_{\text{scale}}, \quad (3.1)$$

where P is the measured output power, V_{out} is the output voltage, and P_{scale} is the full power scale the detector is set to. The two detectors are calibrated relative to each other by using the Toptica to scan across the full range of available wavelengths and measuring the power with each detector.

The lensed fibre tip must be precisely aligned and rotationally adjusted relative to the chip to ensure that light is launched into the desired optical mode of the waveguide. The fibre is positioned in a V-groove holder mounted on a six-axis stage, allowing precise control of its orientation. The pitch, yaw, and roll adjustments are set and locked once the stage is levelled to achieve a perfectly flat configuration. The schematic for the fibre alignment process is shown in Fig. 3.1 (b). To align the polarization, the tip of the lensed fibre is positioned close to a free-space linear polarizer set to 90° , while the in-line HWP is fixed at 0° . The laser is transmitted through the fibre and directed towards the linear polarizer with a silicon photodiode detector (Thorlabs S121C) attached immediately behind it. When the fibre is correctly oriented, the cross-polarized light is extinguished to approximately -20 dB. Throughout the alignment process, care is taken to avoid twisting or pinching the fibre, as this can increase the polarization sensitivity to external disturbances. The alignment procedure is performed for both fibres to ensure consistent coupling conditions across the system.

After the alignment of both fibres, the system throughput is evaluated by positioning the fibre tips in close proximity for optimal coupling and scanning across the full wavelength range. Measurements are first taken with the HWP set to 0° to characterize the TE mode, and subsequently at 45° for the TM mode. For each configuration, the power transmitted through the PBS is recorded to determine the relative contributions of the TE and TM components. When the fibres are correctly aligned, the extinction ratio should remain approximately -20 dB. However, wavelength-dependent oscillations observed in the TE and TM power traces indicate a rotation between the polarization axes of the fibres, resulting in an elliptically polarized beam at the PBS output. The resulting data files for each launch mode contains the wavelength, power from detector 1 (TE mode), and power from detector 2 (TM mode).

For the setup in the cryostat, the procedure is similar, except the cryostat compatible fibre holder does not have a fibre rotator. The fibre rotation is aligned right outside the cryostat as in the photo shown in Fig. 3.1 (b). Inside the cryostat, each fibre is placed on a stack of three piezos allowing for XYZ motion. The sample sits on a pillar. The sample pillar and two piezo stacks sit atop a stage with two piezos, allowing for X and Y movement relative to an objective attached to Z piezo (see. Fig. 3.2). The lensed fibres are connected inside the cryostat to a continuous fibre feedthrough. With the fibres connected,

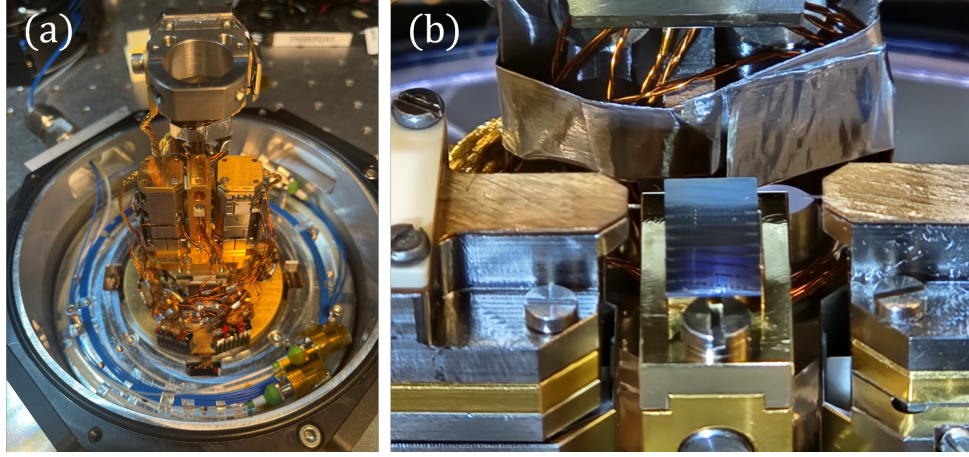


Figure 3.2: Fibre-coupled system inside a cryostat. (a) The opened cryostat without a sample or fibre holders placed on the piezo stages. The fibre fed into the cryostat are coiled along the base of the 40 K stage. (b) Fibre holders with the lensed fibres are installed on either sides of a PIC.

the fibre alignment to the correct orientation and system loss measurement are performed just outside the cryostat. The rotationally aligned fibres are placed on the piezo stacks within the 4 K region, and the extra lengths of the fibre are secured onto the base of the 40 K stage using tape (see Fig. 3.2 (a)).

It was mentioned earlier that the detectors have been measured relative to each other. The detector calibration file is used to adjust the power from detector 2 for each wavelength. The resulting adjusted powers represent the power throughout of the system (i.e., system loss). With this, we can assume an additional power loss is due to the device under test (DUT). It is important to use the same laser power output setting for subsequent measurements corresponding to the calibration file.

With the DUT inserted between the fibre tips, the fibres are brought close to the waveguide of interest, first with the fibre connected to the pump laser then with the collection fibre. The XYZ stages are used to optimize the power through the waveguide. Again, a measurement is taken by launching into each mode (TE and TM) of the waveguide, and the collected transmission passed through a PBS then two detectors.

The measured voltages are converted into a power using Eq. 3.1. Since the Toptica's software records wavelengths to 0.001 pm with potentially many repeats when rounded to 1 pm, in the processing of the data, the wavelengths are rounded to 1 pm and the average power of the repeats is used. For each launch mode, the wavelengths for the DUT transmission file and calibration file are matched, with the non-matching wavelengths discarded. From there, the calibrated transmission at each wavelength is obtained by

$$P_{\text{cal}}(\text{dB}) = 10 \log_{10} \frac{P_{\text{DUT}}}{P_{\text{system}}} \quad (3.2)$$

where P_{cal} is the calibrated power in dB from the DUT, P_{DUT} is the detected transmitted

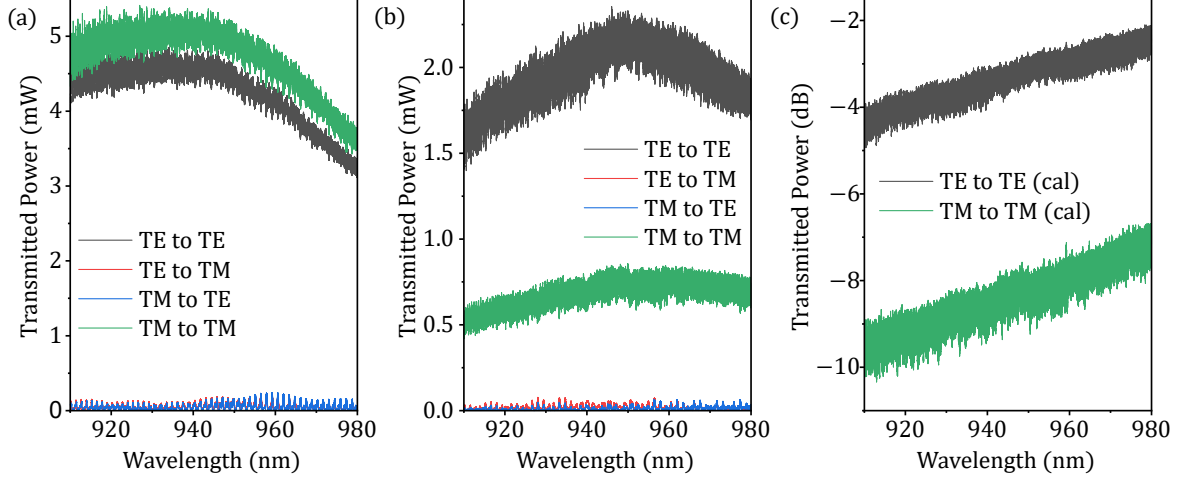


Figure 3.3: (a) Transmission through the system by coupling lensed fibre to lensed fibre. (b) Transmission through the system with device inserted. (c) Calibrated transmitted power in dB through the system only due to the device. Not shown here is the calibrated transmission launching in one polarization and collecting in the perpendicular polarization since the transmitted power for both was so low, the data becomes very noisy after calibration.

power of the DUT in milliwatts, and P_{system} is the transmitted power through the system in milliwatts when coupling the fibres tip to tip.

3.1.2 Loss measurements

A chip containing waveguides matching the dimensions of the other devices of various lengths is included in each foundry run (see Fig. 3.4 (a)). The transmission through each waveguide is used to calculate the loss as a function of the waveguide lengths (i.e., propagation loss) and loss associated with coupling light from a fibre into a waveguide at the edge of the chip (i.e., facet loss).

The data in Fig. 3.3 (c) is smoothed by fitting with a fifth-order polynomial to extract a loss for each corresponding wavelength. A fifth-order polynomial was chosen to obtain acceptable fits across multiple waveguides. The extracted loss at each wavelength as a function the waveguide length (see Fig. 3.4 (b)) is then fitted with a linear function. The slope gives the propagation loss, and the y-intercept gives the loss from the coupling of the lensed fibres to the waveguide facets. Since there are two facets in this case, the loss per facet (i.e., facet loss) is the y-intercept divided by two. It should be noted that the facet loss is dependent on the type of fibre used. For example, for two fibres with the same core diameter, a cleaved fibre placed perfectly flat against the facet will have a lower facet loss than a lensed fibre. Although advantageous for fixed alignment, due to this direct contact with the chip edge, there is a higher chance of damaging both the fibre and the facet. Lensed fibres are preferred for these sets of measurements due to ease of alignment owing to their longer working distance. An optical image of a lensed fibre coupled to a waveguide

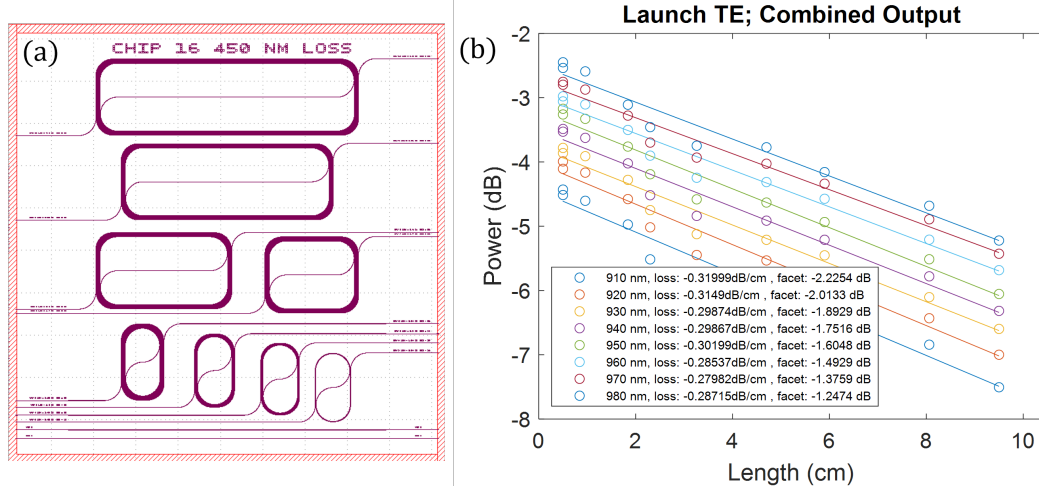


Figure 3.4: (a) Layout of a typical loss chip. The chip shown here is 5 mm by 5 mm. (b) Transmittance as a function of the waveguide length.

is shown in Fig. 3.1 (c). A gap between the fibre tip and waveguide can be clearly seen.

In later chapters, the propagation and facet loss are used for estimating the efficiency of the hybrid integrated single-photon source.

3.2 Nanowire Array

The steps for growing a nanowire array were presented in Ch. 2.2.1, in which an electron beam is used to pattern the array for nanowire growth. A typical finished nanowire growth chip is shown in Fig. 3.5 (a). Each array uses a different dose in $C/\mu\text{m}^2$ which directly affects the size of the hole from which the nanowire will grow. The higher the dose, the larger the diameter of the hole and the resulting nanowire. Since all the nanowires are grown in the CBE simultaneously using the same amount of material, the wider the nanowire, the shorter it will be. This is well illustrated in the “organ pipe” pattern shown in Fig. 3.5 (b).

As the quantum dots in the nanowire core increase in diameter, the less lateral confinement there will be, and the smaller the spacing between the s-, p-, d- orbitals. This is observed when looking at the nanowires pitched $0.4\mu\text{m}$ apart. With the nanowires grown so densely, it is possible to excite and collect from multiple NW-QDs at once using the setup later described in Ch. 3.4.2. Fig. 3.5 (c) shows PL spectra from nanowires grown with a dose of $27886 C/\mu\text{m}^2$ (red) and $36251 C/\mu\text{m}^2$ (blue). The nanowires in both arrays were pumped above-band at the same power. It can be seen that spacing in terms of wavelength between the s-, p-, and d- shells are closer in the higher dose than the lower.

The arrays on the left side of the chip contain nanowires pitched $7.5\mu\text{m}$ apart, and on the right are nanowires pitched $22.5\mu\text{m}$ apart. A magnified image of the arrays on the left side of the chip is shown in Fig. 3.5 (d). Each array is well labelled with the pitch, repeat, and dose. Each row corresponds to a different dose that increases as you go up ranging

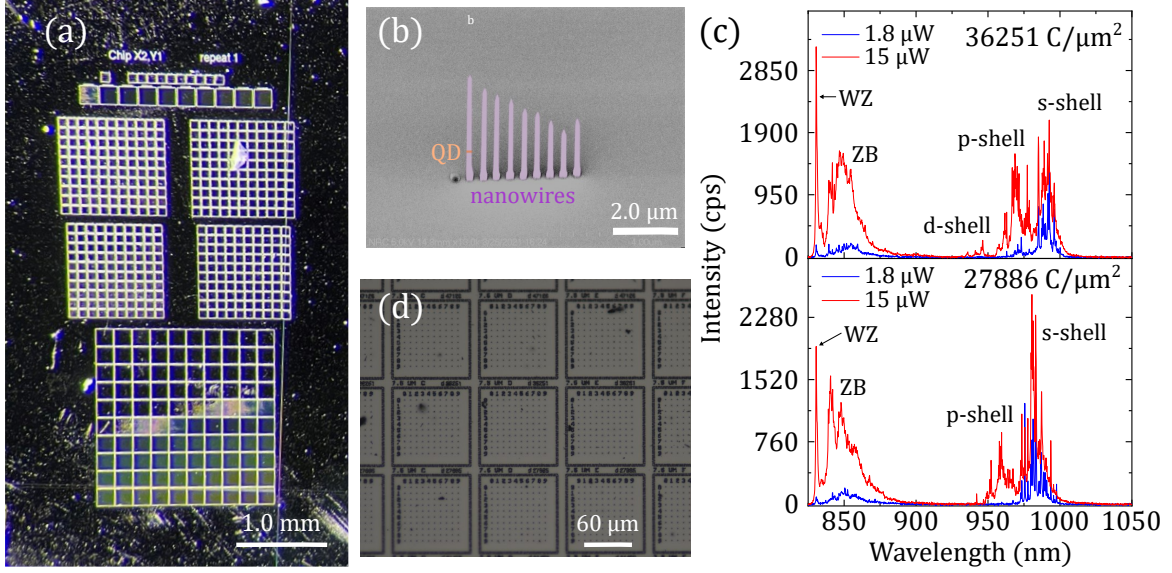


Figure 3.5: (a) Optical image of a nanowire growth chip. (b) SEM image of nine nanowires resulting from increasing electron beam doses from left to right. Openings made at the lowest dose does not always yield a properly grown nanowire. (c) Photoluminescence spectra obtained from nanowires pitched $0.4 \mu\text{m}$ apart from two different doses. (d) Optical image of multiple arrays of nanowires.

from $16500 \text{ C}/\mu\text{m}^2$ to $174974 \text{ C}/\mu\text{m}^2$. Each column is a repeat of the pattern using the same dose labelled from A to K.

For the nanowires presented in Ch. 2, the nanowires are pitched $7.5 \mu\text{m}$ apart. This spacing ensures collection from one NW-QD at a time when using a 100x objective with a numerical aperture (NA) of 0.81. Each of the arrays have nanowires arranged ten by ten with columns and rows labelled 0-9 (see Fig. 3.5 (d)).

The labelling of the arrays and nanowires promotes ease of relocating a specific nanowire. As mentioned before, the nanowires in the arrays on the right side of the chip are pitched $22.5 \mu\text{m}$ apart. This larger spacing is ideal for the pick-and-place method described in the next session, Ch. 3.3, as it allows for a probe tip to move freely within the region without damaging nearby nanowires.

3.3 Hybrid-integration: pick-and-place method

As mentioned in Ch. 2.2.1, the single-photon sources used in this thesis are InAsP quantum dots embedded within nanowires grown in-house at the NRC. This gives us the unique advantage of tailoring the geometry of the nanowires for optimal coupling to Si_3N_4 waveguides.

The motivation behind the pick-and-place method is to enable the deterministic positioning of individual nanowire devices onto Si_3N_4 waveguides [11]. This approach allows

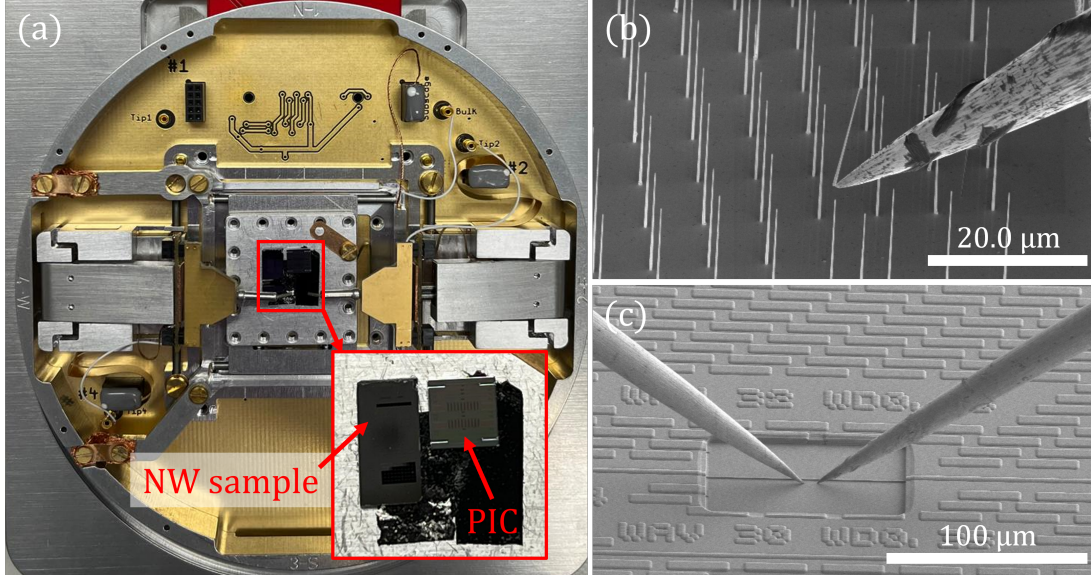


Figure 3.6: Scanning electron microscope-based nanomanipulator. (a) Photograph of the nanomanipulator stage. Inset shows the nanowire growth chip and PIC chip for transfer. (b) SEM image of the probe tip picking up a nanowire. (c) SEM image of the nanowire being transferred to the waveguide.

us to integrate the single-photon emission from InAsP–InP nanowires with the low-loss guiding properties of Si_3N_4 waveguides, creating a scalable and stable on-chip quantum light source.

The integration is performed using a SEM-based nanomanipulator equipped with two independently controlled probes and a central stage that holds both the growth substrate and the Si_3N_4 chip. The setup is shown in Fig. 3.6 (a). The probes can move in three dimensions with nanometre precision through piezoelectric actuators. A target nanowire is first identified within the array on its growth substrate, detached from the growth substrate at its base, and attracted to the probe tip via electrostatic forces (Van der Waals) generated in the SEM environment (i.e., the nanowire is broken off and stuck to the probe tip) as shown in the snapshot in Fig. 3.6 (b). The nanowire and probe tip are then lifted up and moved across to the PIC chip. A second probe tip is used to assist in the positioning and releasing of the nanowire onto the desired location on or beside the SiN waveguide as shown in Fig. 3.6 (c). The nanowire is placed within an exposed region of the waveguide where the protective SiO_2 layer has been etched away. For some chips (like the one used in Ch. 5), the etch depth of the SiO_2 is chosen to optimize coupling of a nanowire placed beside a waveguide rather than on top. This was done to maximize optical mode overlap and ensure efficient evanescent coupling while also making the transfer process easier. Once positioned, the nanowires adhere securely to the waveguide surface without the need for adhesives or further processing.

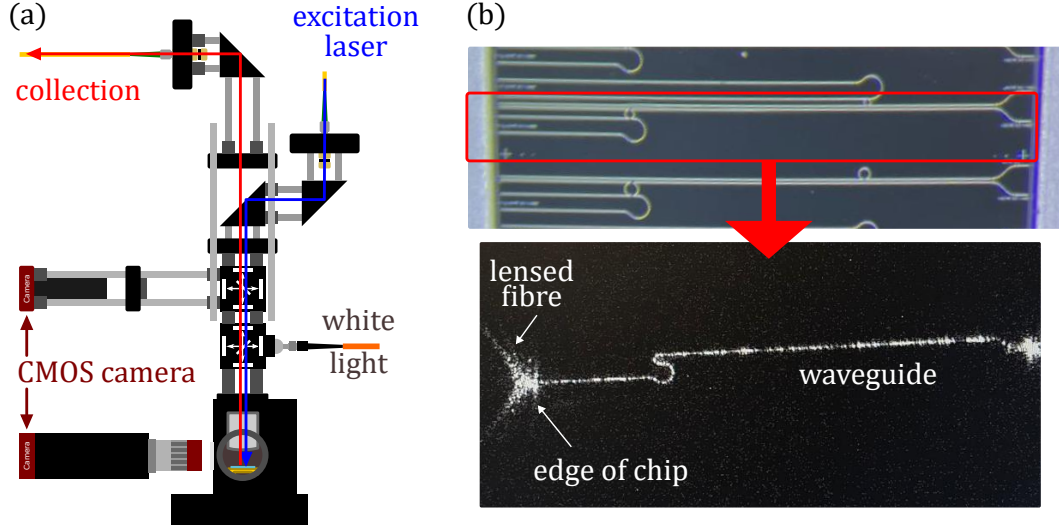


Figure 3.7: (a) Schematic for a free-space pump and collect setup from a cryostat. (b) (top) Optical image of a chip containing a complex waveguide structure. (bottom) Laser light scattering as it travels through a waveguide. The chip is 5 mm wide.

3.4 Optical characterization

This section describes the experimental techniques used to characterize the optical properties of the InAsP quantum dots embedded within InP nanowires after integration with Si_3N_4 photonic circuits. In each case, the chip is cooled to 4 K in a fibre-coupled, closed-cycle helium-cooled cryostat (Attodry 800). These low temperatures provide a more stable environment for the QD with reduced noise in the system. Such an environment enables the QD to emit cleaner emission lines while preserving the coherence of the quantum states for longer. Higher temperatures decrease the coherence due to increasing interactions with phonons. With even higher temperatures, interactions with phonons also result in carriers no longer being confined in the dot.

A 100x objective is used inside the cryostat to assist in fibre alignment, but can also be used as a confocal microscope as in the free-space pump-and-collect setup shown in Fig. 3.7 (a). A monochrome complementary metal-oxide-semiconductor (CMOS) camera (Thorlabs CS165MU) is used for imaging of the chip surface. Scattering of the laser light travelling down the waveguide is used as a visual aid for finding the proper height needed when coupling light from the fibre into the waveguide. An example of laser light scattering through a curved section of waveguide is shown in Fig. 3.7 (b).

The different excitonic complexes of the quantum dot are analysed through different photoluminescence spectroscopy techniques, time-resolved photoluminescence, single-photon purity, and indistinguishability measurements.

3.4.1 Single-photon Detectors

The two types of single-photon detectors used in this thesis include single-photon avalanche diodes (SPADs) and superconducting nanowire single-photon detectors (SNSPDs). Both types convert detected light into an electrical signal that is sent to a counting card (Swabian TimeTagger) to time tag and tally up the number of detection events within a bin of time (e.g., 100 ps). Both are click-no-click detectors unable to resolve the number of photons that have impinged on the detector at a single point in time. The speed and time performance of a single-photon detector is quantified by the uncertainty in the detection time of an arriving photon (timing jitter) and the time duration following a photon detection event for which the detector needs to reset before sensing another photon (dead time). The dead time limits the maximum number of photons that can be detected in a given period of time. Each detector is also quantified by its efficiency in detecting photons in terms of the probability of detecting a single photon (device efficiency) and the rate at which it detects false counts (dark-count rate) [96, 97].

SPADs use an avalanche multiplication process to amplify the signal from a detection event. When a photon impinges on the photosensitive region of the semiconductor diode, through the photovoltaic effect, an electron-hole pair is created. A reverse-bias voltage is applied across the diode that accelerates the electron or hole at high enough energy to create secondary carriers. This in turn causes an avalanche effect that amplifies the photocurrent. The avalanche is stopped with a quenching circuit to reset the detector. Silicon (Si)-based SPADs are used to detect photons in the 260 nm to 1100 nm range with a peak responsivity between 600 nm to 800 nm, and indium gallium arsenide (InGaAs) SPADs operate in the 850 nm to 1700 nm range with a peak responsivity around 1500 nm. In this work, we use Si-based SPADs due to its higher device efficiency in detecting the emission wavelengths produced by our NW-QDs (900 ~ 980 nm). A major advantage of SPADs is that they can be operated at room temperature. However, this leads to a higher rate of dark counts due to thermal and electrical fluctuations in the system [98]. To reduce the dark count rate, SPADs are typically cooled with thermoelectric coolers to temperatures of 210 ~ 250 K [99]. A SPAD typically has a dead time of 10-50 ns, and a timing jitter of 100-200 ps.

SNSPDs consist of a few nanometre thin film of superconducting material operating in the sub 4 K range [97]. It should be noted here that unlike the cylindrical III-V nanowires used throughout this thesis, superconducting nanowires are planar, long meandering wires. There are two leading theories for how an SNSPD works: hotspot model and vortex-crossing model [100, 96], neither of which will be discussed here. In either case, a bias current is applied across the SNSPD. The material stops being superconducting when a photon impinges on the SNSPD, thus resulting in some resistance. As such, a voltage pulse is generated. Key advantages of SNSPDs include their high efficiency, extremely low dark count rates, short timing jitter, fast recovery time, and enhanced sensitivity from the ultraviolet to infrared region as compared to other single-photon detectors such as SPADs, superconducting tunnel junctions (STJ), photo-multiplier tubes (PMT), transition edge sensors (TES), and superconducting hot-electron bolometers (HEB) [97]. A four-channel SNSPD from PhotonSpot with tungsten silicide (WSi₂) superconducting nanowires was

Pump	Collect
F1	F2
F2	F1
F1	F1
F2	F2
F1 or F2	F1 and F2
F1 or F2	objective
F1 or F2	F1, F2, and objective
objective	objective
objective	F1 or F2
objective	F1 and F2
objective	F1, F2, and objective

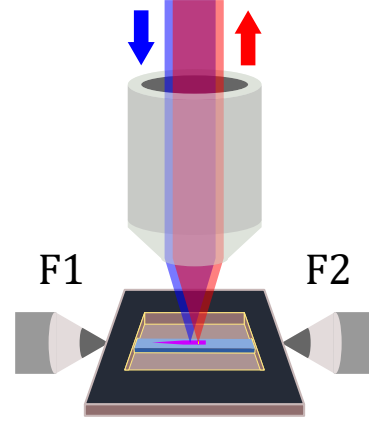


Table 3.1: Pump and collection schemes for a fibre-coupled cryostat system. The figure on the right shows the basic setup with lensed fibres (labelled as F1 and F2) coupled to either sides of the waveguide and an objective above the nanowire.

used in this work.

Specific details on the detector used are given in the text of the individual papers presented in this thesis.

3.4.2 Photoluminescence

Photoluminescence (PL) spectroscopy is employed to measure the emission spectra from an optically excited quantum dot using either CW or pulsed lasers. The laser wavelength is chosen depending on the excitation scheme (i.e., above-band, quasi-resonant, resonant).

With a fibre-coupled system, there are multiple possible pump-collection configurations as listed in Table 3.1 which lists different combinations of free-space and fibre-coupled setups depending on the application. The basic setup for the different schemes is shown beside the table. It is possible to collect from more than one port simultaneously depending on the application. Many of the setups mentioned will be used in Ch. 5 with the exception of collection from two fibres and an objective simultaneously.

One advantage of having separate pump and collection ports when measuring the photoluminescence is that both sides can be optimized independently, whereas in a confocal microscope setup, if the pump and collection beams are of differing wavelengths, they will have different focuses. As a result, if the pump power is actually set too high, as the focus gets closer to the optimal distance for the laser's wavelength, the QD may already become over-saturated despite the collection focus not being close to optimal, thus the overall intensities in the QD emission spectrum being lower. In other words, without decreasing the pump power to a power that will not saturate the s-shell of the QD, finding the optimal focus for pumping and collecting the QD emission can prove difficult. This problem is avoided in the fibre setup where, after rough alignment with a pump laser of a similar wavelength to the expected emission of the s-shell, the QD emission from the collection

	Pylon:100BR_eXcelon	Pixis 100
pixel array	1340×100	1340×100
pixel size	$20 \mu\text{m} \times 20 \mu\text{m}$	$20 \mu\text{m} \times 20 \mu\text{m}$
cooling	-120°C	-75°

	SpectroPro HRS-500
focal length	0.5 m
grating	150 g/mm or 1200 g/mm

Table 3.2: Specifications of the spectrometers – above: cameras; below: monochrometers

fibre can be optimized first on the charge-coupled device (CCD), then the pump fibre alignment can be optimized by looking at the intensity of the wurtzite emission (832 nm) which does not saturate.

In order to pump and collect from the same fibre, either a beam splitter or circulator can be used. It should be noted that in the setup where pump and collection is from the same fibre port, the same problem faced by a confocal setup may occur. However, we typically strive to have the fibre located at a focus to optimize QD collection by aligning as before. In such a case, more laser power may be needed to saturate the QD. This is generally not a major concern since power-dependent measurements are presented relative to the saturation power of the emission line being studied.

Diffraction grating spectrometer

A diffraction grating spectrometer can be used to measure the emission spectrum of a quantum dot. The system can be measured using a combination of free-space and fibre-coupled techniques. However, only the fibre-coupled system will be detailed in this section. On the pump side, the setup is generally the same as discussed above in Ch. 3.1.2 except with a different pump laser depending on the needs of the application. To roughly align the fibres to the waveguide initially, the setup for the collection side is the same as in Ch. 3.1. After rough alignment is achieved, the output fibre is connected to a fibre-coupled grating spectrometer (1200 g/mm) equipped with a silicon-based CCD camera with the specifications listed in Table 3.2. Fig. 3.8 (a) shows the schematic for a diffraction grating spectrometer.

As light enters the spectrometer slit, it diverges and is collimated by a parabolic mirror that directs the beam toward the diffraction grating. The grating disperses the incident polychromatic light such that each wavelength is reflected at a different angle, allowing spectral separation based on the grating equation

$$n\lambda = d(\sin \theta_i + \sin \theta_d), \quad (3.3)$$

where n is the order of the diffraction, λ is the diffracted wavelength, d is the distance between successive grooves, θ_i is the angle of incidence, and θ_d is the angle of diffraction.

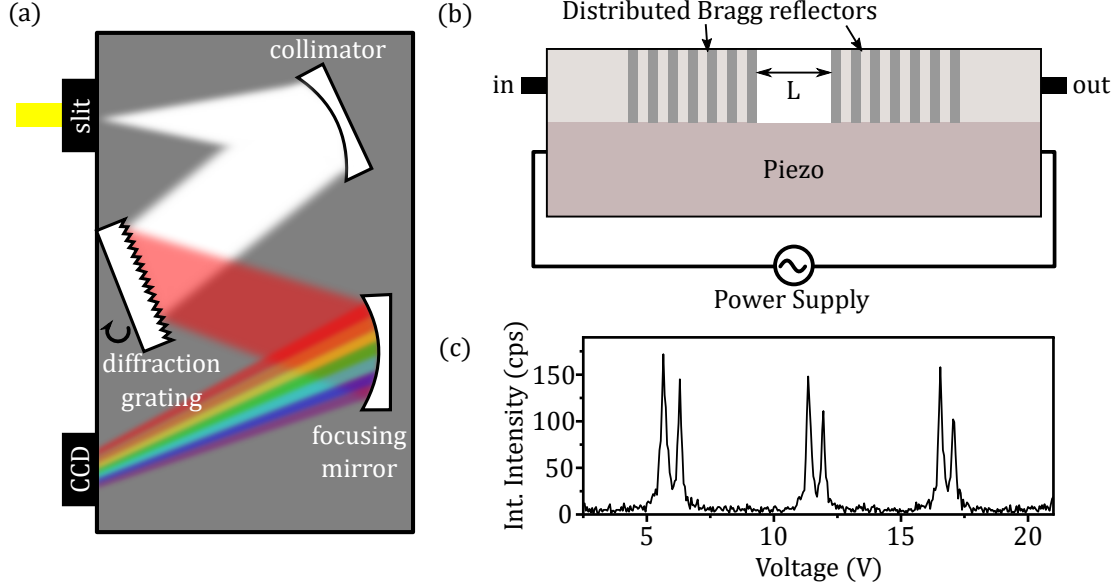


Figure 3.8: (a) Schematic for a diffraction grating spectrometer. (b) Schematic for a high-resolution PL setup using a Fabry-Pérot interferometer. (c) A sample scan for a X photon spanning three repeats (two free spectral ranges (FSRs)).

Another parabolic mirror is used to focus the dispersed light onto the CCD detector. The vertical pixels of the CCD are typically binned when measuring the intensity counts at each wavelength.

Typically for the 1200 g/mm grating centred at 950 nm using the CCDs listed above, the resolution is 0.0552 nm per pixel. For the 150 g/mm grating, the resolution is 0.5199 nm. Using the 1200 g/mm grating, a typical quantum dot emission line spans three pixels which is not enough to resolve the FSS of some excitonic complexes.

High-resolution photoluminescence

To resolve the fine-structure splitting and measure the intrinsic linewidths of the quantum dot emission beyond the resolution limits of a grating spectrometer, high-resolution PL spectroscopy was performed using a fibre-coupled Fabry-Pérot interferometer (i.e., etalon). The schematic of the setup is shown in Fig. 3.8 (b).

A narrow-band tunable filter with a bandwidth of 0.01 nm is used to first spectrally filter the emission to isolate the excitonic transition of interest. The filtered light is then directed through a voltage-controlled Fabry-Pérot consisting of two fibres with mirrors on a piezo with a bandwidth of 200 MHz and a FSR of 40.75 GHz. The cavity length of the etalon is tuned via a piezoelectric element, allowing precise control of the resonance condition. As the cavity length is varied, the transmitted intensity is recorded using a single-photon detector. For the experiments used in this thesis, this would either be a SPAD or SNSPD. The resulting intensity versus voltage profile is reconstructed into a high-resolution emission spectrum (shown earlier in Fig. 2.9 and Fig. 2.10(b)).

To convert from voltage to frequency (GHz), the following equation is used

$$\nu = \frac{\Delta\nu_{\text{FSR}}}{\Delta V_{\text{FSR}}} \times V, \quad (3.4)$$

where V is the voltage setting of the etalon, ν is the corresponding frequency, $\Delta\nu_{\text{FSR}} = 40.75$ GHz is the FSR of the etalon in terms of frequency, and ΔV_{FSR} is the measured FSR in volts. The spectrum in terms of voltage is obtained by scanning the filtered emission across a voltage range showing three repeats as shown in Fig. 3.8 (c). The average of the two peak-to-peak voltages is used as the FSR in Eq. 3.4.

Note, in later chapters showing high-resolution PL spectra, only the centre repeat is shown. For these sets of experiments, the purpose is to measure the linewidth and FSS if present. As such, the frequencies or energies are adjusted relative to the centre of the peak. The conversion from gigahertz (GHz) to microelectron-volts (μeV) is given by the following equation

$$E(\mu\text{eV}) = \frac{h(\nu(\text{GHz}) \cdot 10^9)}{e \cdot 10^{-6}}, \quad (3.5)$$

where Planck's constant $h \approx 6.626 \times 10^{-34}$ J.s and $e \approx 1.602 \times 10^{-19}$ C. The equation can also be rewritten as $E(\mu\text{eV}) \approx 4.1357 \cdot \nu(\text{GHz})$.

The line shapes are fit using Lorentzian or Voigt profiles depending on the degree of homogeneous and inhomogeneous broadening observed. The Lorentzian component corresponds to homogeneous broadening of the emission, whilst the Gaussian component reflects environmental charge noise and spectral diffusion (inhomogeneous broadening). This will be further discussed in Ch. 4. This high-resolution technique allows a direct comparison between measured linewidths and lifetime-limited values calculated from time-resolved photoluminescence measurements.

Power-dependent photoluminescence

Power-dependent PL measures how the intensity and spectral features of the photoluminescence emission vary as a function of the excitation pump power applied to the system. This dependence reveals information about the physical processes governing carrier generation, recombination, and saturation dynamics within the emitter.

To perform a power-dependence PL measurement, the pump power applied to the system is varied by changing the amount of attenuation on the input laser power. In this work, the laser power is controlled using a fibre-coupled, voltage-controlled, electronic variable optical attenuator (VOA). The output power is directed through a beam splitter with one arm going to a power meter for monitoring, and the other arm being sent to the rest of the system.

By varying the excitation power, changes in the quantum dot's emission intensities can be probed as each complex reaches saturation (P_{sat}). This is where all available states are occupied and further excitation does not increase emission intensities. For quantum dots, this method is particularly useful because different excitonic complexes (neutral exciton

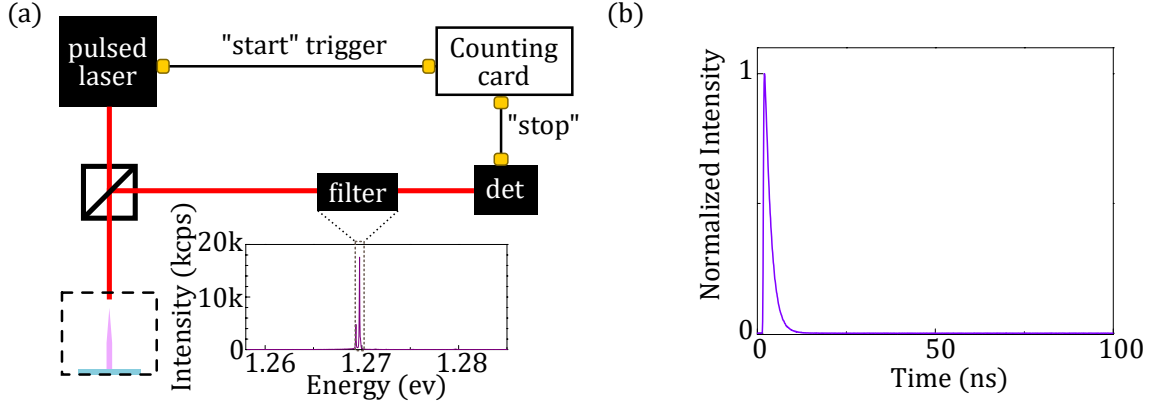


Figure 3.9: (a) Schematic for a lifetime measurement using a counting card to measure the time between a “start” signal from a pulsed laser and a “stop” signal from a photon detection event. (b) Typical histogram produced in a lifetime measurement of an exciton.

X , charged exciton X^- , and biexciton XX photons) show distinct power dependencies, allowing them to be identified and characterized.

For each excitation power, the PL spectrum is recorded. The integrated intensity of each emission line is plotted as a function of excitation power on a log-log scale. As mentioned in Ch. 2.2.5, extracting the slope (n) of the resulting curve can be useful for identifying the excitonic complex.

The saturation curve of the neutral exciton photons (X) is expected to yield $n \approx 1$, meaning its intensity increases proportionally with excitation power. The biexciton photons (XX) is expected to yield $n \approx 2$ because two electron-hole pairs must be present simultaneously. The charged exciton photons (X^-) often falls between these regimes and is expected to have a slope of $n \approx 1.5$ due to the presence of an additional carrier [71].

3.4.3 Time-resolved photoluminescence

Time-resolved photoluminescence (TRPL) is an experimental technique used to measure the temporal dynamics and radiative recombination processes of a quantum emitter. It provides direct access to the excitonic lifetime which is the time it takes for an excited state to recombine and emit a photon after optical excitation. This parameter is crucial for understanding the efficiency and recombination mechanisms of a quantum emitter. This value ultimately determines the repetition rate at which a single-photon source can be operated.

TRPL measurements are commonly carried out using time-correlated single-photon counting. A counting card (Swabian Instruments TimeTagger 20 or TimeTagger Ultra) is used to measure the timing between clicks of two channels: a trigger from the laser and click from the single-photon detection event.

In the setup shown in Fig. 3.9 (a), a pulsed laser is used to excite the QD. There are two main pulsed lasers used in our lab: 1. a pulsed diode laser that operates at 670 nm

with a dial to control the repetition rate at 80 MHz, 40 MHz, 20 MHz, 10 MHz, and 5 MHz, and 2. a Ti:Sapphire ultrafast laser (Mai Tai HP) with a tuneable wavelength output range of 690~1040 nm operating at 80 MHz. Ideally, the repetition rate is chosen such that the QD has enough time to fully decay between sequential pulses. The output signal from the laser is directly connected to the TimeTagger via a coaxial cable to function as the “start” trigger.

To measure the lifetime of a particular excitonic state, the emission is first filtered for the peak of interest. A photon detection event from either an SPAD or SNSPD sends an electrical signal to the TimeTagger telling it to “stop” the timer. Repeating this process over many excitation cycles builds up a histogram of photon arrival times, τ , and builds a picture of the radiative lifetime. A histogram of the times between clicks is constructed as shown in Fig. 3.9 (b).

The curve in Fig. 3.9 (b) can be modelled by [16]

$$I(\tau) \propto e^{-(\tau-\tau_0)/\tau_r} \left[\operatorname{erf} \left(\frac{\tau - \tau_0}{\sqrt{2}\sigma} \right) + 1 \right], \quad (3.6)$$

where τ_0 is the delay offset, τ_r is the radiative decay time. The error function with a rise time of $\sqrt{2}\sigma$ originates from the instrument response function. In the majority of this thesis, we typically do not include the modelling of the rise time. For the X photons, the lifetime may exhibit a double-exponential decay due to the presence of a dark state. This is further discussed in more detail in Ch. 5.

3.4.4 Second-order correlations measurements

For the measurement of an emitter’s single-photon purity, the emission from the source is separated into two streams and sent to two detectors connected to a counting card. The counting card measures the timing difference, τ , between when a photon is detected at a first detector then at a second detector. The two streams can either come from a beam splitter or from two different collection ports on the fibre-coupled hybrid device as listed in Table 3.1. A second-order correlation measurement performed on the same emission complex is referred to as an autocorrelation.

The schematic for a single-port autocorrelation measurement is shown in Fig. 3.10 (a). The emission from the QD is filtered for the wavelength of interest. Then the filtered emission is sent to a beam-splitter, and the outputs directed to two detectors connected to a counting card. The counting card measures the timing delay, τ , between detection events (i.e., “clicks”) across the two channels. The time differences are accumulated into a histogram, where all clicks are considered both as “start” and “stop” clicks for both positive and negative time differences [101]. An anti-bunching dip at $\tau = 0$ with $g^{(2)}(\tau = 0) < 0.5$ would indicate a single-photon source as two simultaneous detection events cannot occur, otherwise there must have been two or more photons impinging on the beam-splitter (BS).

If two collection ports from the emitter are used as shown in Fig. 3.10 (b), then a beam-splitter is not necessary. Instead the directional emission from the source acts as the

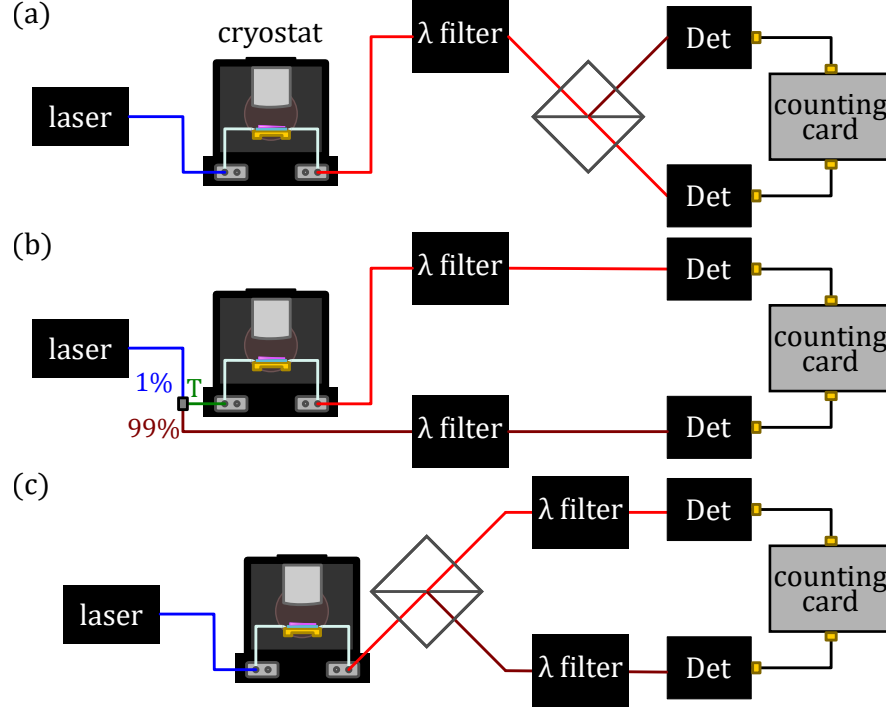


Figure 3.10: (a) Schematic for a single-port autocorrelation $g^{(2)}(\tau)$ measurement setup. (b) Schematic for a double-port second-order correlation setup. For an autocorrelation, the filters are set to the same wavelength. For a cross-correlation, the filters are set to the corresponding wavelengths of each complex being measured. (c) Schematic for a single-port cross-correlation setup. The red, blue, and maroon lines represent optical fibre connections, and the black lines with yellow SMA connectors represent electrical connections.

beam-splitter itself. In this case, for an autocorrelation $g^{(2)}(\tau)$, both filters in each arm are set to filter for the same complex. This method allows for the recuperation of photons that would have otherwise been lost emitting in the opposite direction. The uses for this setup is discussed further in Ch. 5 and Ch. 6.

A cross-correlation $g^{(2)}(\tau)$ measurement can be used to study the temporal relationship between photons emitted from distinct excitonic complexes within a single quantum dot. It provides insight into how two emission lines are related, and whether their emission events are independent or sequentially emitted. Cross-correlation measurements can be particularly useful for verifying the identity of each complex [30, 65]. For example, a strong correlation between the XX and X photons indicates the presence of the biexciton–exciton cascade.

In the setup for a cross-correlation measurement, two filters are needed to filter for each complex under study. Either the setups in Fig. 3.10 (b) or (c) can be used. For these set of measurements, it is important to keep track of which channel is set as the “start” and which is the “stop.” As in the example of the biexciton–exciton cascade, when the “start” is set to the XX photon and the “stop” on the X photon, a strong bunching at $\tau > 0$ indicates that right after the XX photon is emitted, there is a high probability of detecting

the X photon shortly after. In the $\tau < 0$ domain, we see that right after the X is detected, it is unlikely for the XX to be detected shortly after.

3.4.5 Indistinguishability

First-order correlation measurements

A first-order correlation measurement, $g^{(1)}(\tau)$, is used to characterize the coherence properties of a light source. It measures how well the amplitude and phase of a light wave at one time or location is correlated with those at another time or location, and thus directly relate to the temporal coherence and spectral purity of the emitted photons. Measuring the coherence time is essential for understanding indistinguishability of the single-photon source.

There are two typical setups used to measure $g^{(1)}(\tau)$: a Michelson-interferometer or a Mach-Zehnder interferometer (MZI). In both cases, we look at the interference of the light at a beam splitter as the beam recombines. A relative phase shift between the arms can lead to either constructive or destructive interference. Since the experiments in this thesis are all fibre-based, only the MZI will be described here. Fig. 3.11 (a) shows a basic schematic for an MZI featuring two beam-splitters, a phase modulator, and a delay stage. A more detailed setup can be found in Ch. 5. A filtered emission line is first passed through a beam-splitter. In one arm, a phase modulator is used to rapidly change the phase of the light going into the second beam-splitter. For a coherence measurement, the other arm needs to have the same temporal length as the first arm.

To know if the two arms of the Mach-Zehnder interferometer (MZI) are temporally balanced at the second beam-splitter, a pulsed laser can be used in place of the quantum dot. Like in a TRPL setup, the laser output signal is used as the trigger, and the output from one of the arms of the second beam-splitter is used as the click. The resulting histogram looks similar to that of a lifetime measurement.

To balance the two arms for the setup shown in Fig. 3.11(a), the laser is connected to one of the inputs of the first beam splitter. Neither the wavelength filter nor cryostat are used during calibration since only the laser light is being detected. The rest of the setup can remain the same.

However, in the case where the first beam splitter is replaced with the bidirectional coupling setup as in Fig. 3.10(b), both output paths constitute as part of the two arms of the MZI. The distance travelled by the photons through the Si_3N_4 waveguides, lensed fibres, cryostat feedthrough fibres, filters, and other components on each path arm all contribute to the total length for each arm of the MZI before reaching the second beam splitter. To measure the path difference between the two, the QD needs to be excited with a pulsed laser. In this case, the emission from the QD needs to be filtered for both paths.

If the two arms are temporally balanced, the two peaks should perfectly overlap. Otherwise, two peaks can be seen as in the example in Fig. 3.11 (b). The difference in the delay between the two peaks indicate how much longer one arm is relative to the other. Given

that light requires approximately 5 ns to travel through a 1 m fibre, adding that length to one arm enables observation of the relative shift in the peaks, thereby indicating which arm is shorter. In the case of Fig. 3.11 (b), adding a 2 m length of fibre to the shorter arm moved the larger peak (in purple) over by 11 ns (red). Since the length of one arm is unchanged, the smaller peak remained at the same time across the two measurements (circled in blue).

A motorized delay stage in one arm can help correct for a small temporal difference as long as it is within the delay stage's range of travel. For the experiments in this thesis, a delay stage with 1.2 ns of travel was used.

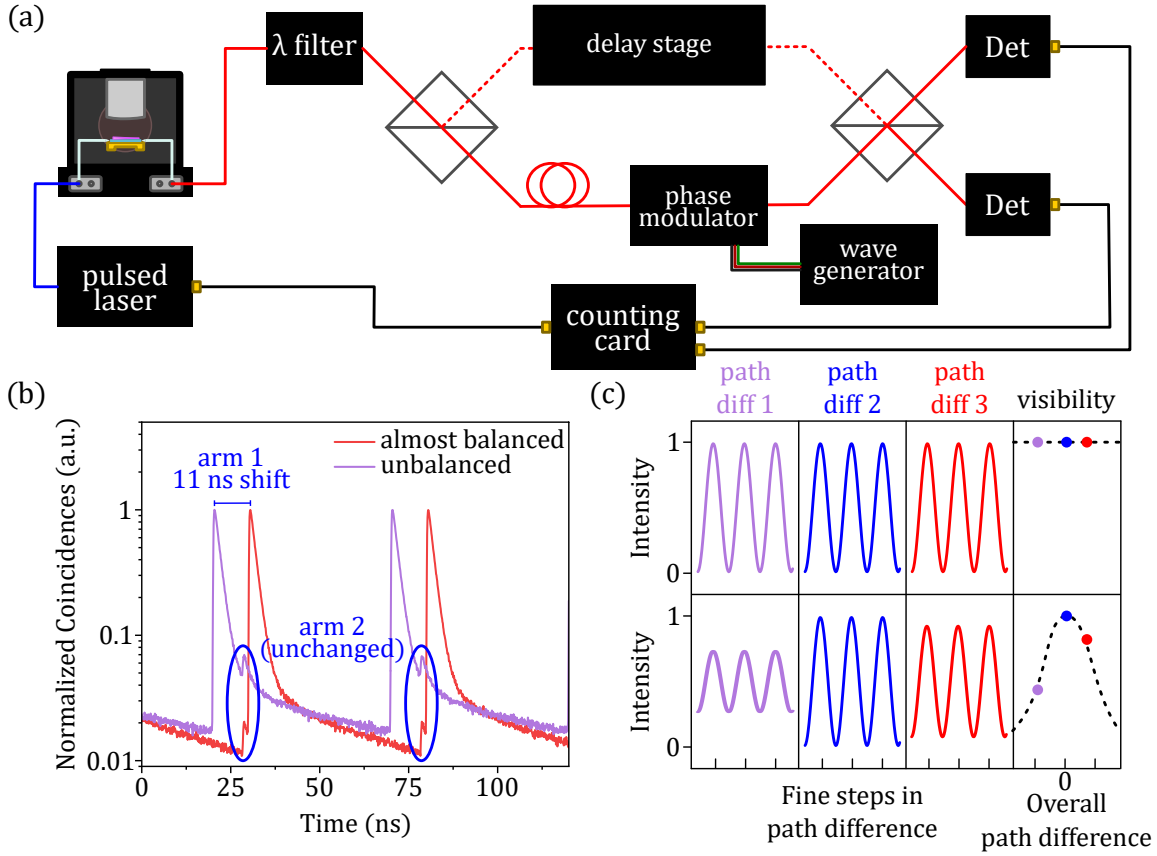


Figure 3.11: (a) Schematic for a Mach-Zehnder interferometer used to measure the coherence of a light source. The red and blue lines represent optical fibre connections, and the black lines with yellow SMA connectors represent electrical connections. (b) Example timing-calibration measurement showing unbalanced arms in an MZI setup. A ~ 0.2 m fibre difference placed in one arm resulted in the 11 ns shift in the peak. (c) Count rate as a function of path difference. (above) A source with a long coherence time such as a laser. (below) A source showing smaller amplitudes in the visibility fringes depending on the path length difference. Visibility (dashed black line) is plotted in the last column.

Interference fringes such as the one shown in Fig. 3.11 (c) can be seen when plotting the count rate detected as a function of time. The fringe visibility at a given delay is obtained

using the phase modulator in one arm of the MZI and is calculated using [9]

$$V = \frac{I_{\max} - I_{\min}}{I_{\max} + I_{\min}} = |g^{(1)}(\tau)|, \quad (3.7)$$

where $I_{\max}$ is the average maximum detected counts (crests) and $I_{\min}$ is the average minimum detected counts (trough). The fringes may also be modelled with a sine wave to better estimate the fringe visibility. As the temporal length between the two arms increases, the fringe visibility decreases. This decay in the fringe visibility can be used to extract the coherence time, T_2 .

For a coherent light source such as a the Toptica CTL with a narrow linewidth (10 kHz), the coherence time should be on the order of milliseconds. As such, the fringe visibility should remain $V \approx 1$ across the entire temporal length of the delay stage. The system is calibrated by measuring V at every 50 ps along the delay stage's range. The resulting curve produced is fitted to smooth the data and used to calibrate subsequent measurements such that

$$V_{\text{cal}}(t) = V_{\text{measured}}(t)/V_{\text{laser}}(t), \quad (3.8)$$

where $V_{\text{cal}}(t)$ is the calibrated fringe visibility at the temporal delay setting, t , of the delay stage, $V_{\text{measured}}(t)$ is the measured fringe visibility, and $V_{\text{laser}}(t)$ is the fringe visibility of the laser at delay setting t . More detail on how to model the fit for extracting T_2 will be discussed in Ch. 4 and Ch. 5.

Two-photon interference

Indistinguishability refers to how identical two photons are in all measurable degrees of freedom which include spectral frequency, polarization, spatial mode, temporal profile, and phase. For a QD in a fibre-based system passing through a narrow-band filter, the spatial mode is in the HE_{11} mode of the fibre. The temporal aspect can be controlled using the Hong-Ou-Mandel (HOM) effect where two identical photons incident on a BS will leave from the same port [102]. To ensure two sequentially emitted photons are incident on the BS simultaneously, a MZI setup is used to measure the two-photon interference (TPI). The temporal delay or path length difference between the two arms is precisely matched to the repetition rate of the pulsed laser used to excite the QD. A CW laser can also be used to get a continuous stream of photons incident on the second BS. Due to spectral broadening, there is an uncertainty in the spectral frequency of the emitted photons even with the use of a narrow pass filter. The polarization of the emitted photons impinging on the second BS can be controlled to measure the degree of indistinguishability of the emitted photons by changing the polarization of the QD emission with a HWP. A basic schematic for the setup is shown in Fig. 3.12 (a).

The output of the beam splitter is sent to two detectors connected to a counting card. Separate $g^{(2)}(\tau)$ measurements are taken with the HWP set to 0° (parallel, $g_{\parallel}^{(2)}(\tau)$) and 45° (perpendicular, $g_{\perp}^{(2)}(\tau)$).

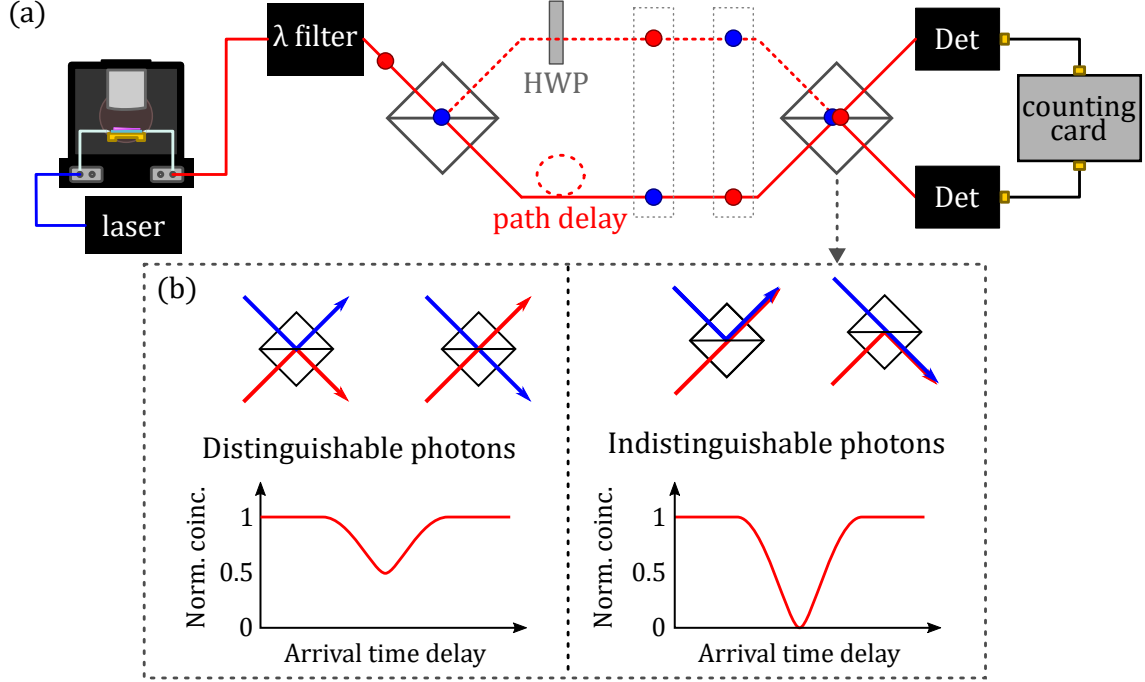


Figure 3.12: (a) Schematic for a two-photon interference measurement setup. Red and blue lines represent optical fibre connections, and black lines represent electrical connections. (b) Possible outcomes at the second beam-splitter due to the Hong-Ou-Mandel effect.

Fig. 3.12 (b) shows the possible outcomes when the two photons arrive at the second beam-splitter. Due to the HOM effect, if the two photons impinging on the BS are indistinguishable (i.e., HWP at 0°), the photons will leave from the same port. This translates to seeing a dip at $g_{\parallel}^{(2)}(\tau=0)=0$. Conversely, if the two photons are completely distinguishable (i.e., HWP at 45°), then the dip goes to $g_{\perp}^{(2)}(\tau=0)=0.5$. This is because half the time, the two photons are leaving from the same port, whilst the other half of the time, they leave from separate ports.

The TPI visibility is calculated using

$$V_{\text{TPI}} = \frac{g_{\perp}^{(2)}(0) - g_{\parallel}^{(2)}(0)}{g_{\perp}^{(2)}(0)}. \quad (3.9)$$

For perfectly indistinguishable photons, $V = 1$, while $V = 0$ indicates complete distinguishability.

Chapter 4

III-V Nanowire Quantum Dot on Basic Photonic Integrated Circuit

This chapter expands on work done during my MSc [18] in which nanowire-quantum dots (NW-QDs) were integrated onto straight photonic waveguides and tested for their photoluminescence and single-photon purity. Here, additional experiments were performed on the sample to measure the source's coherence time and indistinguishability. The details related to the setup of the measurement was discussed in Ch. 3.4.

4.1 Article: On-chip indistinguishable photons using III-V nanowire/SiN hybrid integration

This chapter includes the following publication and its corresponding supplementary materials:

Reprinted with permission from **Yeung, E.**, Northeast, D. B., Jin, J., Laferrière, P., Korzusinski, M., Poole, P. J., Williams, R. L., & Dalacu, D. (2023). On-chip indistinguishable photons using III-V nanowire/SiN hybrid integration. *Phys. Rev. B*, 108(19), 195417. <https://doi.org/10.1103/PhysRevB.108.195417>. Copyright 2023 American Physical Society (APS).

4.1.1 Contribution of co-authors

E. Yeung conceptualized and performed all the experiments. D. B. Northeast designed the chip layout. E. Yeung wrote the code related to controlling equipment, automating the experiments, and compiling the data. E. Yeung and D. B. Northeast performed the pick-and-placement of nanowires onto the PIC, with D. B. Northeast acting as a mentor. J. Jin helped with the first iteration of the Hong-Ou-Mandel (HOM) interferometer setup. The final HOM setup was designed by E. Yeung, D. B. Northeast, and D. Dalacu. D. B.

Northeast calculated the required fibre lengths for both the HOM and coherence setups and spliced the fibres.

E. Yeung and D. Dalacu compiled and analyzed the data. M. Korkusinski helped with the theoretical interpretation of the data.

E. Yeung, D. B. Northeast, D. Dalacu, P. J. Poole, and R. L. Williams were all in active discussion related to the contents and editorial feedback of the paper. E. Yeung was the lead in writing the paper.

D. Dalacu supervised the research.

On-chip indistinguishable photons using III-V nanowire/SiN hybrid integration

Edith Yeung,^{1,2} David B. Northeast,¹ Jeongwan Jin,¹ Patrick Laferrière,¹ Marek Korkusinski,^{1,2} Philip J. Poole ¹,
Robin L. Williams,¹ and Dan Dalacu ^{1,2}

¹National Research Council Canada, Ottawa, Ontario, Canada K1A 0R6

²Department of Physics, University of Ottawa, Ottawa, Ontario, Canada K1N 6N5



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We demonstrate on-chip generation of indistinguishable photons based on a nanowire quantum dot. From a growth substrate containing arrays of positioned-controlled single-dot nanowires, we select a single nanowire which is placed on a SiN waveguide fabricated on a Si-based chip. Coupling of the quantum dot emission to the SiN waveguide is via the evanescent mode in the tapered nanowire. Postselected two-photon interference visibilities using continuous-wave excitation above band and into a p shell of the dot were 100% consistent with a single-photon source having negligible multiphoton emission probability. Visibilities over the entire photon wave packet, measured using pulsed excitation, were reduced by a factor of 5 when exciting quasiresonantly and by a factor of 10 for above-band excitation. The role of excitation timing jitter, spectral diffusion, and pure dephasing in limiting visibilities over the temporal extent of the photon is investigated using additional measurements of the coherence and linewidth of the emitted photons.

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I. INTRODUCTION

High interference visibility between two single photons incident on separate input ports of a 50/50 beam splitter, i.e., the Hong-Ou-Mandel (HOM) effect [1], establishes the indistinguishable nature of the photons, an essential requirement in most photonic quantum technologies [2]. Epitaxial semiconductor quantum dots offer a solid-state solution for deterministically generating indistinguishable photons [3] with state-of-the-art sources demonstrating two-photon interference (TPI) visibilities in excess of 95% [4–6] for pulse separations of over 14 μ s [7] in devices with efficiencies of up to 57% [8]. These sources were designed to emit out of plane, whereas a key advantage of solid-state emitters is the ability to integrate them with on-chip photonic circuitry [9]. An integrated platform, whereby multiple sources generate indistinguishable photons [10,11] propagating within on-chip photonic circuitry [12], is a long-term challenge addressing the scalability requirements of complex quantum processing schemes [2].

Two distinct technologies for generating on-chip indistinguishable photons using quantum dots are currently being pursued. In monolithic approaches [13], the quantum dot is embedded within photonic crystal waveguides [14,15] or suspended nanobeams [16] fabricated from the same III-V material system. Hybrid platforms [17], on the other hand, combine III-V-based quantum dot systems with Si-based integrated circuits in which the dot emission is coupled to waveguide structures using either direct butt-coupling [12] or evanescent fields [18]. Initial experiments [14–16,19] demonstrating the on-chip generation of indistinguishable photons using the above approaches relied on postselected TPI visibilities, where excitation was provided by a continuous-wave (cw) laser. More recently, nonpostselected measurements made using resonant pulsed excitation have shown TPI visibilities of

$V > 90\%$ between sequentially emitted photons separated by a few nanoseconds [20–23] up to almost 800 ns [24].

In this work, we report on a deterministic hybrid technique for generating in-plane indistinguishable photons. We use position-controlled nanowire quantum dots [25] incorporated in photonic nanowire waveguides that are designed for efficient evanescent coupling to SiN waveguides [26]. The approach is similar to previous techniques employing evanescent coupling of III-V structures to underlying Si-based photonic structures [18,27,28], except here the adiabatic mode transfer is not realized through geometries defined by lithography, but rather by a taper introduced in the photonic nanowire during growth [29]. The hybrid construction relies on a pick and place technique, whereby individual nanowires are picked up from the III-V growth substrate and placed on a SiN waveguide fabricated on a Si wafer. The transfer is carried out in a scanning electron microscope (SEM) which provides a placement precision of a few nanometers, i.e., a positioning control sufficient to achieve optimal nanowire-waveguide coupling.

We obtain peak TPI visibilities of $V = 19.2\%$ over the temporal extent of the emitted photons when exciting nonresonantly. To gain insight into the different mechanisms limiting the observation of high visibilities, we perform additional experiments sensitive to decoherence processes in two-level systems: first-order correlation, $g^{(1)}(\tau)$, and high-resolution spectroscopic measurements. These additional measurements point to fluctuations in the charge environment on timescales of < 2 ns as the primary mechanism preventing the observation of high TPI visibilities, i.e., fluctuations expected when employing nonresonant excitation.

II. CHIP-INTEGRATED NANOWIRE SOURCE

The quantum dots used in this work are sections of InAsP incorporated within position-controlled InP nanowires

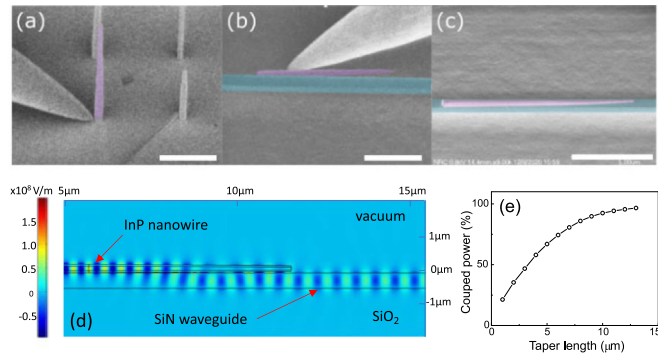


FIG. 1. Pick and place technique: (a) Nanowire (purple) picked up from growth substrate using a nanomanipulator probe. (b) Nanowire placed next to a SiN waveguide (blue). (c) Integrated InP nanowire-SiN ridge waveguide device. Scale bars are 5 μm . (d) Simulation of the electric field E_x of the fundamental waveguide mode propagating through a cross section of the hybrid device. (e) Calculated nanowire to waveguide mode coupling as a function of the nanowire taper length.

grown using gold-catalyzed vapor-liquid-solid epitaxy described in detail in Refs. [30–32]. The nanowires are clad with an InP shell to create a photonic nanowire with a base diameter of 250 nm that supports single-mode waveguiding of 1.3 eV photons. The photonic nanowire is tapered to a tip diameter of 100 nm over a $\sim 15 \mu\text{m}$ length [see Fig. 1(a)] to enable adiabatic mode transfer, discussed below. The low-loss photonic circuitry was fabricated in SiN grown by low-pressure chemical vapor deposition with measured propagation loss of $\sim 0.48 \text{ dB/cm}$ at 965 nm. Waveguide dimensions were 400 nm wide and 485 nm thick with SiO_2 below and above, designed to support a single polarization mode at 1.3 eV. The SiN waveguides were terminated at etched facets of the Si chip where the width was tapered to 250 nm for efficient coupling to single-mode lensed fibers (SMLFs) with measured coupling losses of $\sim 2.15 \text{ dB/facet}$.

The nanowires were transferred to the Si chip prefabricated with the SiN photonic circuitry using an SEM-based nanomanipulator, as shown in Figs. 1(a)–1(c). A single nanowire is picked up from the growth substrate with a tungsten tip controlled by piezomotors and then moved to the Si chip mounted next to the growth substrate. The nanowire is then placed either on or beside a selected SiN waveguide which has been exposed by opening a $50 \times 50 \mu\text{m}^2$ window in the top oxide; see Fig. 2(a). As mentioned above, the nanowires are tapered to promote adiabatic mode transfer from the InP nanowire to the SiN waveguide. In Fig. 1(d), we show a simulation of the evanescent transfer of an appropriately polarized HE_{11} mode in the nanowire to the transverse magnetic mode in the SiN waveguide. From this, we calculate a transfer efficiency in excess of 90% for the nanowire geometry specified above; see Fig. 1(e).

The chip was cooled to 4 K in a fiber-coupled closed-cycle He cryostat equipped with xyz piezostages for fiber-waveguide alignment. The nanowires were optically excited through the waveguide via a SMLF. The emission was collected on the other end of the waveguide through another SMLF and

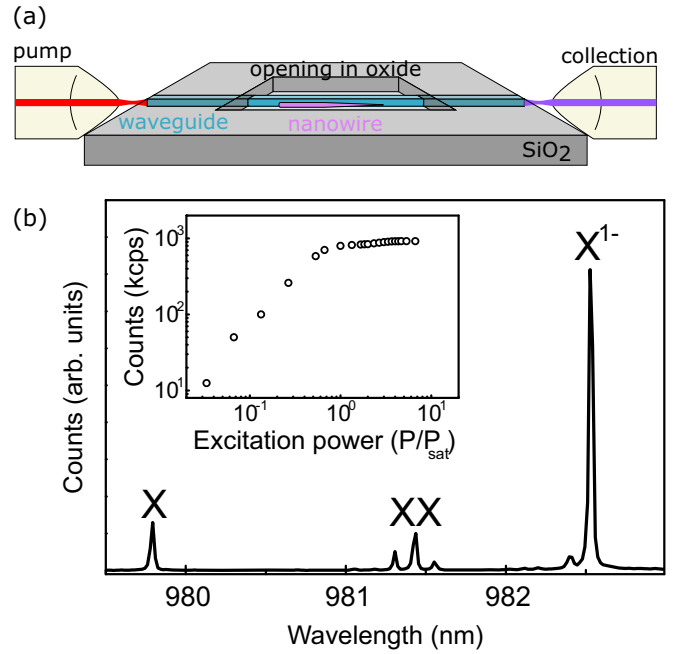


FIG. 2. (a) Schematic of the pump and collection scheme. (b) Photoluminescence spectrum from a hybrid integrated device. Inset: The detected count rates as a function of pump power using pulsed quasiresonant excitation at 80 MHz.

directed to a fiber-coupled grating spectrometer equipped with a nitrogen-cooled charge-coupled device for spectrally resolved measurements or to superconducting nanowire single-photon detectors (SNSPDs, timing jitter 100 ps) via a fiber-coupled tunable filter (bandwidth = 0.1 nm) for measurements on single lines. An *s*-shell photoluminescence (PL) spectrum from the on-chip source is shown in Fig. 2(b) and displays the typical exciton complexes (X , XX , X^{1-}) observed from such nanowire quantum dots [33].

In this article, we focus on the X^{1-} emission. We first determine the efficiency of single-photon generation from the device. For pulsed above-band excitation at 80 MHz, we measure a maximum of 0.919 Mcps at the SNSPD at an excitation power $P = P_{\text{sat}}$ that saturates the transition; see inset in Fig. 2(b). Taking into account the optical throughput of the system (8.1%) and the detector efficiency (88.5%), both measured at 980 nm, we obtained a first-lens count rate of 12.9 Mcps, corresponding to source efficiency of $\eta_s \sim 16.1\%$. To estimate the dot to SiN waveguide coupling, we consider (i) a calculated dot- HE_{11} coupling of $\beta = 95\%$ [34], (ii) a 50% loss of photons for emission directed towards the base of the nanowire, (iii) a 50% loss due to the waveguide design which only supports one of the two polarization modes in the nanowire, and (iv) a 20% loss of photons emitted into phonon sidebands [34] which are filtered out. Taking into account these losses, we obtain an evanescent coupling efficiency of $\eta_c \sim 84.8\%$, slightly lower than both the calculations shown in Fig. 1(e) and our best measured results to date, $\eta_c \sim 93\%$ (see Ref. [35]). The lower-than-predicted values may be associated with emission into other charge complexes [25], e.g., the neutral complexes X and XX , evident in Fig. 2(b).

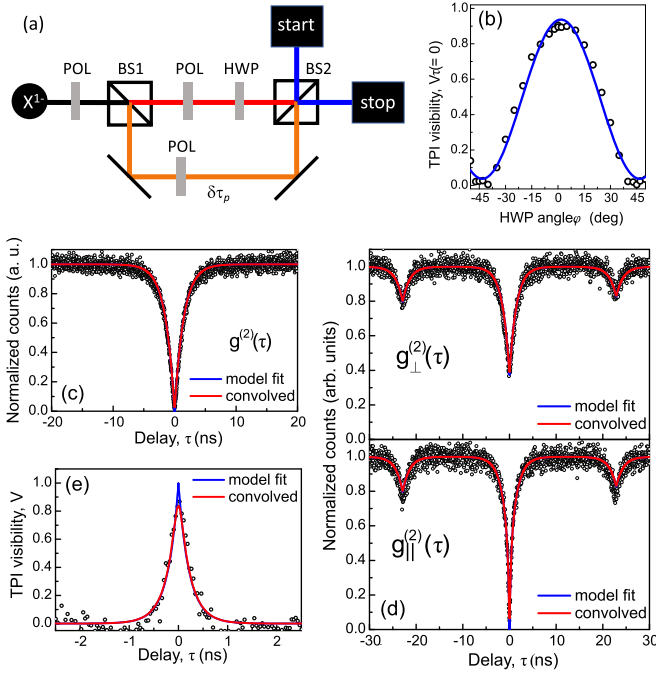


FIG. 3. (a) Schematic of Hong-Ou-Mandel (HOM) setup. (b) Raw TPI visibility at $\tau = 0$ ns as a function of HWP angle ϕ measured using above-band cw excitation at P_{sat} and a $\delta\tau_p = 1.84$ ns delay in one arm of the MZI. (c), (d) Correlation measurements performed with above-band cw excitation at $P_{\text{sat}}/4$ and a $\delta\tau_p = 22.9$ ns delay. (e) Two-photon interference visibility calculated from (d). Black symbols: measured data. Blue curves: model fits described in the text. Red curves: after convolution with detector response.

III. TPI MEASUREMENTS: CW EXCITATION

In this section we describe the TPI measurements performed using cw excitation. This measurement provides information on the interference visibility between sequentially emitted photons where the experiment “selects” those photons that arrive at a beam splitter simultaneously. The extracted visibility is hence a visibility at a delay $\tau = 0$ and an ideal single-photon source, i.e., zero second-order correlations $g^{(2)}(\tau)$ at $\tau = 0$, should be 100%. Extracted values which are lower than this are due to the timing jitter in the detection system. We determine this zero-delay TPI visibility for different experimental conditions, including above-band and below-band excitation and as a function of the delay time between the two emitted photons.

To measure TPI visibilities, the polarization of the X^{1-} photons is first aligned to the slow axis of a polarization-maintaining (PM) fiber using a fiber paddle polarization controller and filtered using the tunable filter to isolate the X^{1-} line. The photons are then input to a PM fiber-based unbalanced Mach-Zehnder interferometer (MZI) equipped with two fiber-coupled free-space nonpolarizing beam splitters (BS1 and BS2), each with a 50:50 nominal splitting ratio; see Fig. 3(a). Two additional polarizers are placed on each arm of the MZI to ensure linear polarized photons incident on BS2. One arm of the MZI includes a half-wave plate (HWP) for controlling the relative polarizations, ϕ , of the photons incident on the input ports of BS2. The delay between the

two arms of the interferometer, $\delta\tau_p$, is adjusted by adding additional fiber to one of the arms. Two SNSPD detectors at the output ports of BS2, labeled “start” and “stop,” together with counting electronics are used to register coincidences.

We first verified the performance of the interferometer by measuring the dependence of the TPI visibility at $\tau = 0$ ns on ϕ , given by $V(0, \phi) = [g^{(2)}(0, \phi) - g^{(2)}(0, 0^\circ)]/g^{(2)}(0, \phi)$. This is plotted in Fig. 3(b) (black symbols), measured using above-band cw excitation and a delay of $\delta\tau_p = 1.84$ ns. We observe an expected oscillatory behavior in the visibility, from $V \sim 0.9$ to $V \sim 0$, as ϕ is varied from 0° to $\pm 45^\circ$. The maximum observed visibility is limited by the detector response, discussed below.

To determine the TPI visibilities as a function of delay τ , coincidence spectra are measured for cross-linear ($\phi = 45^\circ$) and colinear ($\phi = 0^\circ$) polarized photons incident on BS2, $g_{\perp}^{(2)}(\tau)$ and $g_{\parallel}^{(2)}(\tau)$, respectively, and the visibilities are calculated from

$$V(\tau) = [g_{\perp}^{(2)}(\tau) - g_{\parallel}^{(2)}(\tau)]/g_{\perp}^{(2)}(\tau). \quad (1)$$

Typical spectra are shown in Fig. 3(d) (black symbols) for measurements made using above-band cw excitation at $P_{\text{sat}}/4$ and a delay $\delta\tau_p = 22.9$ ns.

To model the coincidence spectra, we examine the four possible path combinations that two photons can take to traverse the MZI and simultaneously arrive at BS2 at $\tau = 0$. Consider first the two cases where the photons are incident on different input ports of BS2 (e.g., the photons travel in different arms of the MZI). If the pair is distinguishable (e.g., in the polarization degree of freedom by setting $\phi = 45^\circ$), each photon is equally likely to be reflected or transmitted, and of the four possible outcomes, only the two where the photons exit different ports will register coincidence counts at zero delay. This is the case in the upper panel of Fig. 3(d), which shows a $g_{\perp}^{(2)}(\tau = 0) \sim 0.5$ relative to the coincidence counts at long delays from uncorrelated photons. If the pair is indistinguishable, these two outcomes will destructively interfere (i.e., the incident pair coalesces, always exiting the same port) so that there are no possible outcomes that can register a coincidence at zero delay. This is the case in the lower panel of Fig. 3(d), which shows $g_{\parallel}^{(2)}(\tau = 0) \sim 0$. The other two cases, where the photons travel in the same arm of the MZI, will both have photons incident on the same port of BS2. Since the photons are generated by the same single-photon source, they cannot arrive simultaneously, and hence will not register zero-delay coincidences.

The absence of simultaneous photons incident on BS1 also eliminates one of the four possible outcomes that would lead to coincidences at $\tau = \pm\delta\tau_p$. This manifests as a reduction in $g^{(2)}(\tau)$ to ~ 0.75 at $\tau = \pm\delta\tau_p$, which is seen in Fig. 3(d).

For 50:50 beam splitters, the behavior above can be modeled by

$$g_{\perp}^{(2)}(\tau) = \frac{1}{2}g^{(2)}(\tau) + \frac{1}{4}[g^{(2)}(\tau - \delta\tau_p) + g^{(2)}(\tau + \delta\tau_p)], \quad (2)$$

$$g_{\parallel}^{(2)}(\tau) = \frac{1}{2}g^{(2)}(\tau) + \frac{1}{4}[g^{(2)}(\tau - \delta\tau_p) + g^{(2)}(\tau + \delta\tau_p)] \times [1 - Fe^{-2|\tau|/\tau_c'}], \quad (3)$$

which describe the cross-polarized, $g_{\perp}^{(2)}(\tau)$, and copolarized, $g_{\parallel}^{(2)}(\tau)$, coincidences, respectively, in terms of the second-order correlation function $g^{(2)}(\tau)$. The first (second) term in Eqs. (2) and (3) accounts for the cases where the photons are incident on the same (different) input port(s) of BS2, while F accounts for the spatial overlap of the photons on BS2 which is assumed to be 100% (i.e., $F = 1$). The timescale τ'_c represents, phenomenologically, the temporal extent over which photons incident on BS2 will coalesce, providing a measure of the probability of coalescence. The functional form with which it is incorporated in Eq. (3) will depend on the mechanism limiting coalescence, e.g., homogeneous versus inhomogeneous broadening. For simplicity, we use an exponential decay (i.e., the visibility is limited by pure dephasing [36]), though we do expect spectral diffusion to play a role due to the nonresonant excitation.

In this work, we use a modified model which explicitly includes the transmission (T_{BS1} , T_{BS2}) and reflection (R_{BS1} , R_{BS2}) coefficients of BS1 and BS2 [37],

$$g_{\perp}^{(2)}(\tau) = 4(T_{BS1}^2 + R_{BS1}^2)R_{BS2}T_{BS2}g^{(2)}(\tau) + 4R_{BS1}T_{BS1}[T_{BS2}^2g^{(2)}(\tau - \delta\tau_p) + R_{BS2}^2g^{(2)}(\tau + \delta\tau_p)], \quad (4)$$

$$g_{\parallel}^{(2)}(\tau) = 4(T_{BS1}^2 + R_{BS1}^2)R_{BS2}T_{BS2}g^{(2)}(\tau) + 4R_{BS1}T_{BS1}[T_{BS2}^2g^{(2)}(\tau - \delta\tau_p) + R_{BS2}^2g^{(2)}(\tau + \delta\tau_p)] \times [1 - Fe^{-|\tau|/\tau'_c}]. \quad (5)$$

The additional terms are required to account for deviations from a perfect, lossless system, i.e., deviations which result in coincidence dip depths that differ from the values described above. For example, if $T_{BS1} \neq R_{BS1}$, then the drop in coincidences at $\tau = \pm\delta\tau_p$ will be less than 25%, and, in the case of distinguishable photons, $g_{\perp}^{(2)}(\tau = 0) < 0.5$. For $T_{BS2} \neq R_{BS2}$, there will be an asymmetry in the depths of the $\tau = \pm\delta\tau_p$ side dips. Although there is little evidence of the latter in the measured spectra (i.e., $T_{BS2} \sim R_{BS2}$), we do observe values of $g_{\perp}^{(2)}(0) < 0.5$ which we associate not necessarily to an unbalanced BS1, but to a difference in the propagation losses between the two arms of the MZI (e.g., from inclusion of the components and accompanying mating sleeves) that result in different count rates incident on the two ports of BS2.

To reproduce the experimental $g_{\perp}^{(2)}(\tau)$ and $g_{\parallel}^{(2)}(\tau)$ using Eqs. (4) and (5), we first measure $g^{(2)}(\tau)$ by detecting coincidences at the two output ports of BS2. The $g^{(2)}(\tau)$ measured under the same operating conditions as the HOM experiment shown in Fig. 3(d) is shown in Fig. 3(c) (black symbols). The response is modeled using $g^{(2)}(\tau) = 1 - e^{-(1/T_1 + R)|\tau|}$ (see the Supplemental Material [38]) where the radiative lifetime $T_1 = 1.75$ ns is independently measured and the excitation rate R is a fit parameter. The blue curve in the figure shows the model fit using $R = 0.1$ ns⁻¹.

This expression of $g^{(2)}(\tau)$ is incorporated in Eqs. (4) and (5), which are then applied to the measured correlations $g_{\perp}^{(2)}(\tau)$ and $g_{\parallel}^{(2)}(\tau)$, with τ'_c and the beam-splitter coefficients as fit parameters. The resulting fits are shown in Fig. 3(d) (blue curves) where, for this particular measurement (above-band

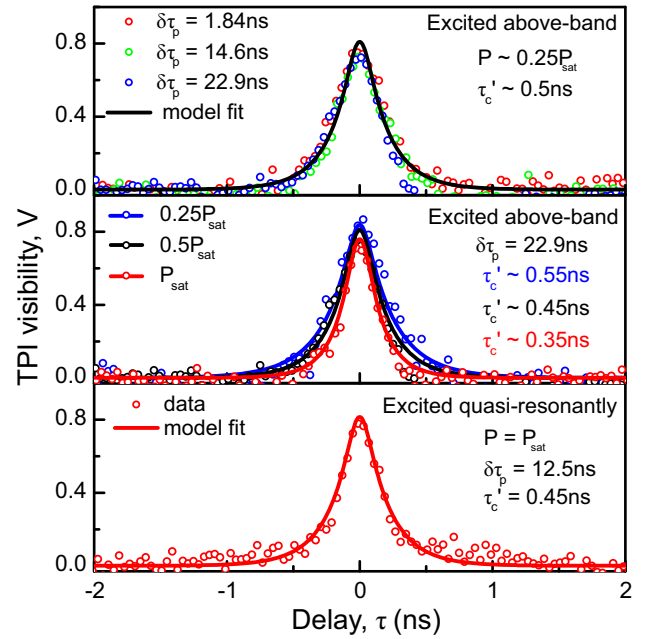


FIG. 4. TPI visibilities obtained under different operating conditions. Upper panel: visibilities using different delays $\delta\tau_p$. Middle panel: visibilities as a function of excitation power. Lower panel: visibility under quasiresonant excitation into a p shell of the dot.

cw excitation at $P_{\text{sat}}/4$ and $\delta\tau_p = 22.9$ ns delay), we obtain a $T_{BS1} : R_{BS1}$ ratio of 0.25:0.75, a $T_{BS2} : R_{BS2}$ ratio of 0.48:0.52, and $\tau'_c = 0.55$ ns.

The visibility obtained from Eq. (1) using the model correlations is shown in Fig. 3(e) (blue curve) and, by definition, predicts a TPI visibility of $V(\tau = 0) = 1$. To match the experiment, we convolve the model correlations with a 100 ps Gaussian detector response function and obtain the red curves in Figs. 3(c)–3(e) which correctly predict the measured raw visibility of $V(\tau = 0) \sim 0.85$.

We have performed the above analysis on measurements taken under different operating conditions to extract τ'_c values. The results are summarized in Fig. 4. For measurements as a function of $\delta\tau_p$ using above-band cw excitation at $P = 0.25P_{\text{sat}}$ (upper panel), we obtained $\tau'_c \sim 0.5$ ns, independent of path delay for $\delta\tau_p = 1.84$ to 22.9 ns. This suggests that the mechanisms limiting the τ'_c values occur on timescales faster than ~ 2 ns. In contrast, from measurements as a function of excitation power (middle panel), we observed a significant increase in τ'_c from 0.35 to 0.55 ns for a fourfold reduction in power. We also observed a moderate increase in τ'_c when we excite quasiresonantly into a p shell of the dot. For measurements at $P = P_{\text{sat}}$, $\tau'_c = 0.45$ ns for p -shell excitation (bottom panel) compared to $\tau'_c = 0.35$ ns for above-band excitation.

Although cw HOM measurements reveal the presence of decoherence mechanisms through the measurement of τ'_c , the extraction of meaningful values requires simultaneous fitting of the three correlations $g^{(2)}(\tau)$, $g_{\perp}^{(2)}(\tau)$, and $g_{\parallel}^{(2)}(\tau)$ measured under the same operating conditions. Without the additional information provided by $g^{(2)}(\tau)$, erroneous values of τ'_c are possible due to the dependence of the antibunching dip in cw measurements on the excitation rate R ; see, for example,

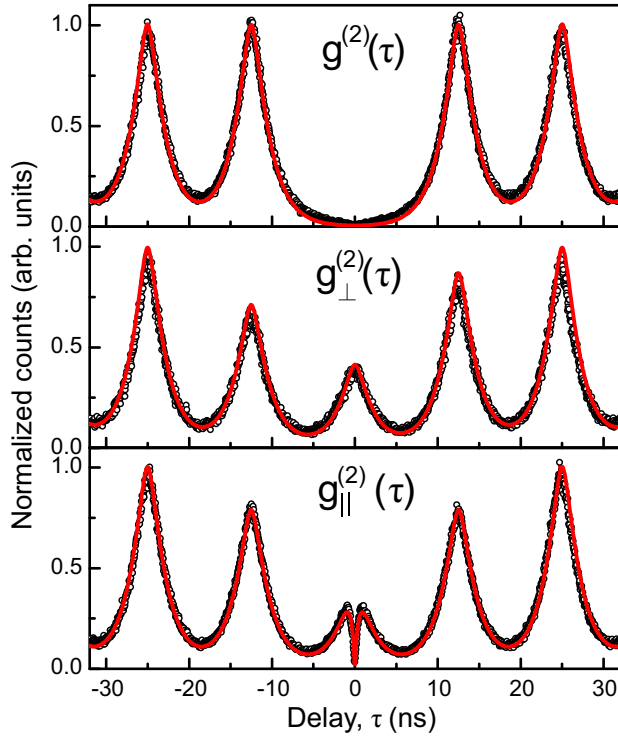


FIG. 5. Correlation measurements for quasiresonant pulsed excitation at $P_{\text{sat}}/4$ (black symbols). Red curves are model fits described in the text.

Ref. [33] and the Supplemental Material [38], Fig. S1. Finally, we note that in all cases, $V(\tau = 0) = 1$ if account is taken of the detector response. This is expected in cw HOM measurements if $g^{(2)}(\tau = 0) = 0$, as is the case here, i.e., $V(\tau = 0)$ is only limited by the detector response [10,37,39].

IV. TPI MEASUREMENTS: PULSED EXCITATION

In this section, we repeat the above measurements but here we excited with a pulsed laser. Using pulsed excitation provides information on the TPI visibilities over the temporal extent of the emitted photons, e.g., nonpostselected. Unlike the cw measurements in the previous section, here we can quantify the probability of emitting identical photons. Excitation using a pulsed source will also be ultimately required for on-demand operation.

Using a tunable pulsed laser (pulse width ~ 2 ps), we measure visibilities as in the previous section, for both above-band and quasiresonant excitation. We use a pulse repetition rate of 80 MHz, i.e., a pulse period of 12.5 ns, and a corresponding delay in one arm of the MZI of $\delta\tau_p = 12.5$ ns. The measured coincidences $g^{(2)}(\tau)$, $g_{\perp}^{(2)}(\tau)$, and $g_{\parallel}^{(2)}(\tau)$ (black symbols) for the case of quasiresonant excitation at $P_{\text{sat}}/4$ are shown in Fig. 5.

Similar to the case of cw excitation for a nominal interferometer (50:50 beam splitters, 100% transmission), if two photons arriving simultaneously at BS2 are distinguishable, we expect a peak centered at zero delay having half the height of the peaks at $\pm 2T$ while the peaks at $\pm T$ should be reduced by 25%. For two perfectly indistinguishable photons arriving

at BS2, the zero-delay peak should be absent. The measured $g_{\perp}^{(2)}(\tau)$ (middle panel in Fig. 5) qualitatively reproduces the behavior expected of impinging distinguishable photons on BS2. For the indistinguishable case, however, the zero-delay peak in the measured $g_{\parallel}^{(2)}(\tau)$ (bottom panel in Fig. 5) is still present, but with a dip that drops to zero coincidences at $\tau = 0$. This behavior is well documented [40–44] and is a consequence of processes that limit coalescence to timescales that are shorter than the temporal extent of the photons, e.g., (i) pure dephasing [45], (ii) spectral diffusion [46], and (iii) excitation timing jitter [47], and thus limit extracted τ'_c to values less than $2T_1$. We note that the last mechanism, i.e., timing jitter, is the primary distinction between the pulsed and cw measurements: it is absent in the latter where the experiment selects only the photons that arrive simultaneously at BS2.

As in the cw case, the curves in Fig. 5 are modeled using Eqs. (4) and (5), but with the latter modified to (see Supplemental Material [38])

$$g_{\parallel}^{(2)}(\tau) = 4(T_{\text{BS1}}^2 + R_{\text{BS1}}^2)R_{\text{BS2}}T_{\text{BS2}}g^{(2)}(\tau) + 4R_{\text{BS1}}T_{\text{BS1}}\{[T_{\text{BS2}}^2g^{(2)}(\tau - \delta\tau_p) + R_{\text{BS2}}^2g^{(2)}(\tau + \delta\tau_p)] - [T_{\text{BS2}}^2g^{(2)}(-\delta\tau_p) + R_{\text{BS2}}^2g^{(2)}(+\delta\tau_p)]Fe^{-2|\tau|/\tau'_c}\}, \quad (6)$$

such that $g_{\parallel}^{(2)}(\tau)$ is simply given by $g_{\perp}^{(2)}(\tau)$ less the exponential term defining the timescale over which the photons coalesce. This is seen clearly for the case of nominal 50:50 beam splitters in which case Eq. (6) reduces to $g_{\perp}^{(2)}(\tau) - 0.5Fe^{-2|\tau|/\tau'_c}$ after setting $g^{(2)}(-\delta\tau_p) = g^{(2)}(+\delta\tau_p) = 1$.

The $g^{(2)}(\tau)$ that is incorporated in Eqs. (4) and (6) is constructed from peaks described by $e^{(-\tau/T_1)}[1 - e^{(-\tau/\tau_c)}] * e^{(\tau/T_1)}[1 - e^{(\tau/\tau_c)}]$, i.e., we self-convolve model fits to the measured time-resolved PL spectra; see the Supplemental Material [38], Fig. S2. This allows us to account for the dependence of the $g^{(2)}(\tau)$ peaks on operating conditions through τ_c , i.e., the timescale associated with preparation of the X^{1-} state. We neglect reexcitation which reduces the single-photon purity to $\sim 98\%$ and assume a $g^{(2)}(\tau) = 0$ in the zero-delay peak. The resulting fits are shown in the figure (red curves) where, as for the cw measurements, we have fit the beam-splitter ratios and τ'_c . We note that for the pulsed measurements, we obtained slightly different beam-splitter ratios for cross- and colinear measurements, which is evident in the data from the pronounced asymmetry between the peaks at $\tau = \pm T$ in the $g_{\perp}^{(2)}(\tau)$ whereas these peaks are symmetric in the $g_{\parallel}^{(2)}(\tau)$ correlations. For the particular measurement shown in Fig. 5, i.e., quasiresonant excitation at $P_{\text{sat}}/4$, we obtain a $T_{\text{BS1}} : R_{\text{BS1}}$ ratio of 0.27:0.73 (0.31:0.69) and $T_{\text{BS2}} : R_{\text{BS2}}$ ratio of 0.35:0.65 (0.5:0.5) for $g_{\perp}^{(2)}(\tau)$ [$g_{\parallel}^{(2)}(\tau)$] and $\tau'_c = 0.95$ ns.

In Fig. 6, we plot a zoom-in of the the zero-delay peaks for both cross- and colinear measurements, $g_{\perp}^{(2)}(\tau)$ and $g_{\parallel}^{(2)}(\tau)$, respectively, for quasiresonant (upper panel) and above-band (lower panel) excitation. The raw visibility over the temporal extent of the emitted photons is determined as in the previous section using Eq. (1), but here we use the integrated coincidence counts over $\tau \pm T/2$. Under quasiresonant

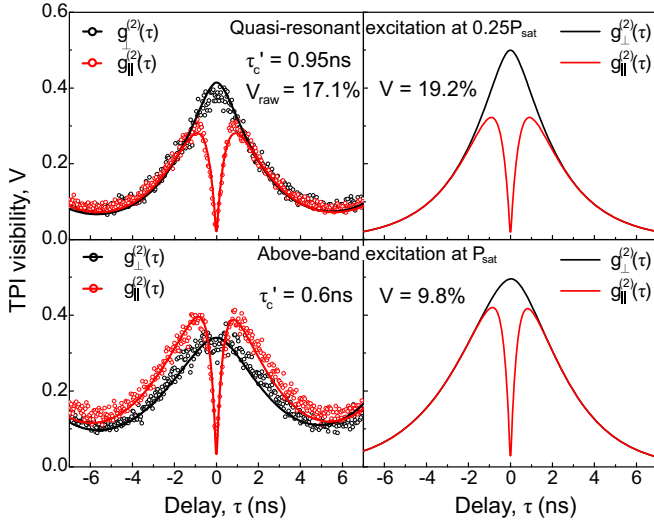


FIG. 6. Zero-delay correlation peaks under quiresonant (upper panels) and above-band (lower panels) excitation. The left panels show the raw data (symbols) and model fits (curves). The right panels show the model calculations with BS1 and BS2 ratios set to 50:50 and contributions from side peaks removed.

excitation at $P = P_{\text{sat}}/4$, we obtain a nonpostselected visibility of $V = 17.1\%$. If the contributions to the coincidence counts from the correlation side peaks are removed (see right panel of Fig. 6), we obtain a corrected visibility (e.g., the visibility that would result using a longer pulse period T) of $V = 19.2\%$. In this plot, we have also corrected for non-nominal values of the beam-splitter ratios for comparison with the above-band measurements, discussed below.

For above-band excitation (lower panel of Fig. 6), the zero-delay peak is significantly broader than that observed when exciting quiresonantly [e.g., compare $g_{\perp}^{(2)}(\tau)$ (black curves)] and the dip at $\tau = 0$ is significantly narrower [compare $g_{\parallel}^{(2)}(\tau)$ (red curves)]. Both broadening of the peak and narrowing of the dip (quantified by the reduced $\tau'_c = 0.6$ ns) are a consequence of the more significant timing jitter [48] present in above-band excitation, i.e., longer timescale associated with the state preparation τ_e , since above-band excitation includes additional processes related to carrier thermalization and capture not present for quiresonant excitation. We note that here we have used a higher excitation power than that used in the quiresonant measurement and this will also contribute to a decrease in τ'_c , as observed in the cw measurements.

We also observe that for the above-band measurements, the zero-delay peaks of cross- and colinear correlations do not overlap at τ values away from $\tau = 0$. This is a consequence of the more significant difference in the beam-splitter ratios between the respective measurements and strongly impacts the calculated visibility using the raw data. Instead, we calculate $g_{\perp}^{(2)}(\tau)$ and $g_{\parallel}^{(2)}(\tau)$ using the fit value of τ'_c , but with 50:50 beam-splitter ratios as above. For the above-band nonpostselected visibility, we obtain $V = 9.2\%$, substantially lower compared to the quiresonant case and consistent with the reduced τ' value.

Comparison with previously reported values is restricted due to the dependence of the measured visibilities on both excitation conditions [42,49] and the delay between interfered photons, $\delta\tau_p$ [50]. For above-band excitation, we refer to the reported values from our previous work [43] where visibilities of $V \sim 5\%$ were measured. There, a temporal delay of $\delta\tau_p = 50$ ns was used and the devices were grown at lower temperature where more severe linewidth broadening is expected; see Ref. [51]. In the case of p -shell excitation with $\delta\tau_p \sim 12.5$ ns, similar visibilities, in the range $V = 0.21$ – 0.59 , have been previously reported [7,50,52], including from on-chip sources [20].

V. COHERENCE MEASUREMENTS

In this section, we determine the coherence properties of the two-level excitonic system. These measurements provide information on the factors (e.g., pure dephasing and charge noise) responsible for the low TPI visibilities observed in the previous section. We show results in the time domain where coherence times τ_c are extracted from single-photon interference visibilities, e.g., $g^{(1)}(\tau)$ measurements. We also compare with results in the frequency domain, where coherence times are extracted from linewidth measurements.

A. Interferometric measurements

For the time-domain measurements, the MZI was balanced (e.g., $\delta\tau_p = 0$) and a motorized fiber-based delay stage with a tuning range of 1.2 ns was added to one arm. The stage was scanned across its full range and, at selected delays $\delta\tau_p$, the fringe visibility was determined using a phase modulator in one arm of the MZI. The fringe visibility as a function of $\delta\tau_p$ extracted from above-band and quiresonant measurements is plotted in Fig. 7.

For a homogeneously broadened transition, i.e., the spectral linewidth of the emitted photons corresponds to the natural linewidth, the line shape is Lorentzian and the visibility is expected to decay exponentially with a time constant T_2 . In the presence of inhomogeneous broadening, the decay will have a Gaussian component [53], T_G , and is more appropriately described by a Voigt profile [43],

$$g^{(1)}(\delta\tau_p) \sim \exp \left[-\frac{\pi}{2} \left(\frac{\delta\tau_p}{T_G} \right)^2 - \frac{|\delta\tau_p|}{T_2} \right]. \quad (7)$$

We model the above-band fringe visibility using Eq. (7) (curves in the upper panel of Fig. 7) to extract T_2 and T_G . To compare with the linewidth measurements below, we also calculated the full width at half maximum (FWHM) of the Voigt profile in the frequency domain given by

$$\delta\omega_V = 0.535\delta\omega_L + \sqrt{0.217\delta\omega_L^2 + \delta\omega_G^2}, \quad (8)$$

where $\delta\omega_L = \frac{1}{\pi T_2}$ is the Lorentzian contribution to the linewidth and $\delta\omega_G = \frac{\sqrt{2\ln 2}}{\sqrt{\pi} T_G}$ is the Gaussian contribution. The extracted Voigt linewidths $\delta\omega_V$ using above-band excitation are plotted in Fig. 9 (filled symbols).

In the case of the quiresonant measurements (lower panel of Fig. 7), instead of a simple decay in the visibility with

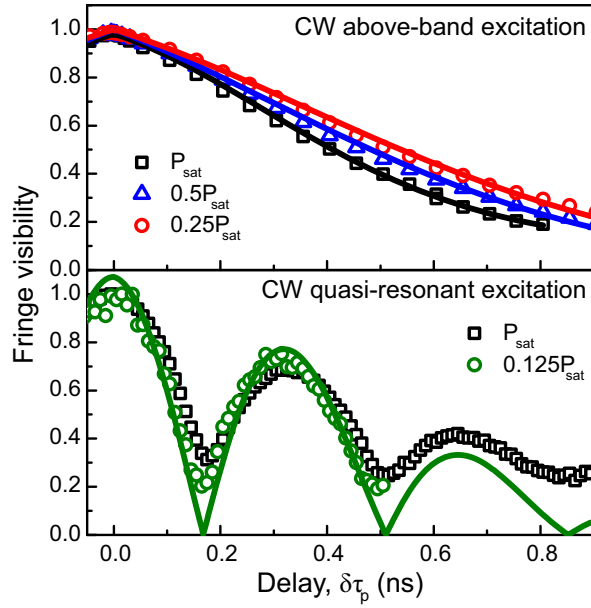


FIG. 7. Fringe visibility extracted from single-photon interference measurements using above-band (upper panel) and quasiresonant (lower panel) excitation. Curves are model fits described in the text.

$\delta\tau_p$, we observe oscillatory behavior indicative of a beating phenomenon. Why this arises is discussed below; here we simply assume the presence of two lines with frequency separation ω_s and identical coherence and model this behavior by multiplying Eq. (7) by the beating term $|\cos(\omega_s\delta\tau_p)|$. We only model the measurement at $P = 0.125P_{\text{sat}}$; at higher excitation powers (e.g., $P = P_{\text{sat}}$), the assumption of identical coherence

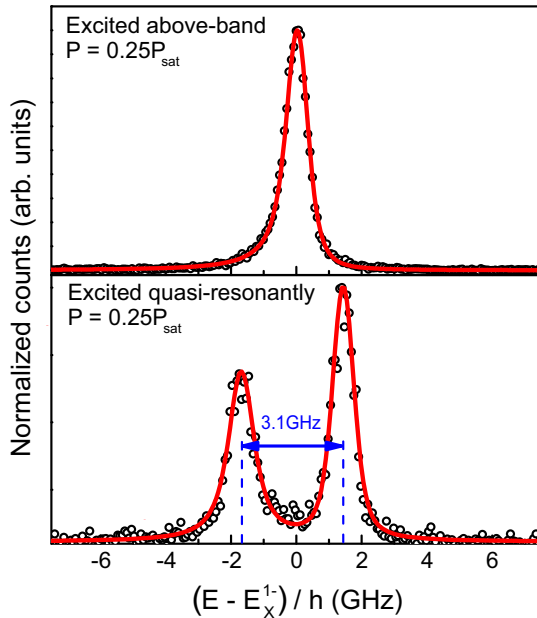


FIG. 8. High-resolution photoluminescence spectrum of the X^{1-} emission for cw above-band excitation (upper panel) and cw quasiresonant excitation (lower panel) at $P_{\text{sat}}/4$. Red curves are Voigt fits to the data.

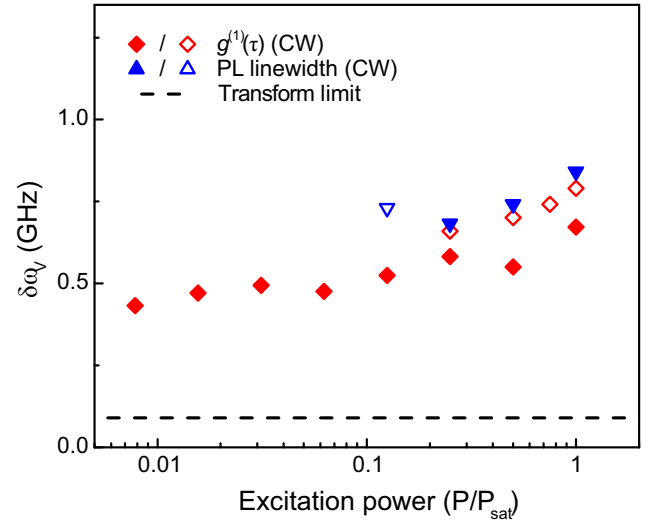


FIG. 9. Voigt linewidths extracted from $g^{(1)}(\tau)$ and high-resolution PL measurements as a function of excitation power. Filled (open) symbols correspond to above-band (quasiresonant) excitation. Dashed horizontal line indicates the transform-limited linewidth, $\delta\omega = 1/2\pi T_1$.

properties likely does not apply, as suggested by the higher minima values in the visibility that are observed in the figure. The extracted $\delta\omega_V$ from Eq. (8) from the quasiresonant pump at $P = 0.125P_{\text{sat}}$ is plotted in Fig. 9 (open symbol).

B. Linewidth measurements

In the frequency domain, we measure the linewidth of the emitted photons using a fiber-based, piezo-controlled Fabry-Perot etalon (bandwidth 250 MHz, free spectral range 40.75 GHz). High-resolution spectra obtained by scanning the etalon through the X^{1-} emission peak are shown in Fig. 8 for cw above-band (upper panel) and quasiresonant (lower panel) excitation at $P_{\text{sat}}/4$. For above-band excitation, we observe a single peak as expected from a singly charged complex, while for quasiresonant excitation, the same emission line is a doublet with a splitting of $\omega_s/\pi = 3.1$ GHz, consistent with the beating frequency observed in the coherence measurements. Given that $g^2(0) \sim 0$, the observation of two peaks suggests either two mutually exclusive excitonic complexes or the same excitonic complex with the electrostatic environment jumping between two possible states.

There are two possible reasons why the observation of either of the two complexes or the same complex in two environments would depend on the excitation energy. First, in quasiresonant excitation, fewer carriers are introduced into the system, resulting in a different Fermi level profile compared to above-band excitation. This may modify the relative intensities of different charge complexes in time-integrated PL spectra, which we have previously observed. Second, there may be a single, defect-related trap sufficiently close to the dot such that a Stark-mediated shift of the X^{1-} emission energy of $\omega_s \sim 3$ GHz will result, depending on the occupation of the trap. In this scenario, for above-band excitation, the trap occupation is constant on the timescale of the measurement (~ 100 ms), whereas for quasiresonant excitation, the

occupation fluctuates on a much faster timescale such that both emission energies are observed in the time-integrated PL.

The magnitude of the observed splitting depends on the quiresonant excitation power, decreasing by ~ 0.4 GHz as the power is increased from $0.25P_{\text{sat}}$ to P_{sat} . Such an excitation power-dependent splitting is consistent with carrier screening of the Stark field, but difficult to explain in the case of two distinct complexes. We therefore attribute the two peaks observed when exciting quiresonantly to the same excitonic complex X^{1-} .

High-resolution spectra were measured as a function of excitation power, deconvolved from the etalon response and fit using a Voigt line shape. We note that the line shapes obtained using above-band excitation present a slight asymmetry when excited at higher powers and these were fit with an asymmetric Voigt function. The origin of this asymmetry is unclear, but is typical of the nanowire quantum dot system [51]. The total linewidth $\delta\omega_V$ from the fits is plotted in Fig. 9 for both the above-band (filled) and quiresonant (open) excitation. For clarity, only measurements of one of the two peaks observed with quiresonant excitation are included (the power-dependent linewidths of the second peak are similar).

Comparison of the time- and frequency-domain measurements reveal similar linewidths which do not significantly depend on the excitation energy: above band (filled symbols) versus quiresonant (open symbols). The absence of any narrowing when exciting below band is surprising: in fact, we observe a small increase in the measured linewidth for quiresonant excitation. This suggests that the mechanisms responsible for the excess broadening are unrelated to the much higher density of excess carriers or phonons introduced with above-band excitation.

In all cases, we observe a decrease in excess broadening at lower excitation powers with a minimum extracted linewidth of $4X$ the transform limit (dashed line in Fig. 9). The reduction in linewidth with excitation power is expected from a reduction in both phonon-related pure dephasing and charge noise-related spectral diffusion. We do not attempt to quantify these two contributions in terms of the Lorentzian versus Gaussian contributions to the measured linewidth: although the total Voigt linewidth is considered accurate, the Lorentzian:Gaussian ratio likely has a large uncertainty.

VI. DISCUSSION

To compare the coherence measurements with the HOM results, we use T_2 and T_G extracted from the $g^{(1)}(\tau)$ and linewidth measurements to calculate a “coherence time” τ_c given by [46]

$$\tau_c = -\frac{T_G^2}{\pi T_2} + \sqrt{\left(\frac{T_G^2}{\pi T_2}\right)^2 + \frac{2T_G^2}{\pi}}. \quad (9)$$

The calculated values of τ_c from the three experiments are summarized in Fig. 10. All measurements give τ_c values of 0.5 ± 0.1 ns at $P = P_{\text{sat}}$, which increase as the excitation power is reduced, consistent with the linewidth

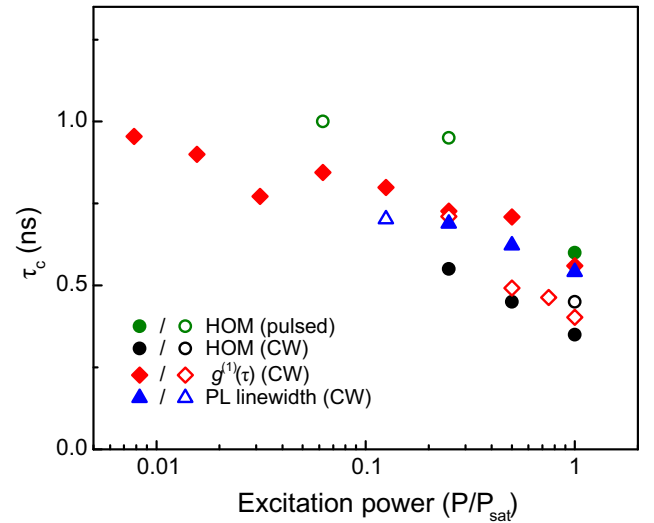


FIG. 10. Coherence times extracted from HOM, $g^{(1)}(\tau)$ and high-resolution PL measurements as a function of excitation power. Filled (open) symbols correspond to above-band (quiresonant) excitation.

measurements in Fig. 9, but still well below the transform limit of $2T_1$.

In the following, we consider the distinct nature of each experiment to extract information on the nature of the mechanisms responsible for $\tau_c < 2T_1$. We first consider the different timescales associated with each experiment. In the HOM measurement, different photons are interfered on a timescale of 2–23 ns, while in the coherence measurement, interference is between the same photon and the timescale is 1.2 ns, but each point in the fringe visibility trace is a statistical average over ~ 100 ms. In the linewidth measurements, on the other hand, there is no interference, only a measure of the spectral purity over timescales in the seconds. The consistency in the extracted coherence times irrespective of the experiment thus suggests that the mechanism(s) responsible for reducing values of τ_c below the transform limit of $2T_1$ occur on a timescale of < 2 ns and there are no other mechanisms on longer timescales for at least seconds.

Next we note the absence of any significant improvement in τ_c when switching from above-band to quiresonant excitation, which is in stark contrast to other studies [42]. This suggests that phonon dephasing is not the primary mechanism limiting the coherence times since the phonon occupation, by necessity, will be higher for above-band excitation [51], which leaves spectral diffusion as the likely source, meaning the charge environment is equally stable, regardless of excitation mode.

Finally, we note that all the measurements, with the exception of the pulsed HOM experiments, are independent of excitation jitter. And yet it is these measurements that produced the highest values of τ_c (compare, for example, pulsed versus cw HOM in Fig. 10). This suggests that either timing jitter is relatively unimportant or that pulsed excitation produces a more stable charge environment, thus compensating for any reduction in τ_c due to timing jitter.

VII. CONCLUSION

In summary, we have demonstrated the generation of indistinguishable photons on-chip based on the hybrid integration of position-controlled single quantum dot nanowires and silicon-based photonic integrated circuitry [26]. Nonpostselected measurements over the photon lifetime revealed coalescence of only a small fraction of the emitted photons, $\sim 20\%$. The TPI visibility was limited by the excess broadening that arises when exciting nonresonantly. Higher visibilities are anticipated with coherent excitation, as demonstrated in other quantum dot systems [4–6].

The described integration approach provides a route to developing a platform whereby multiple sources of indistinguishable photons can be selectively incorporated on-chip, a long-term goal of future quantum technologies. As a final note, in this work we investigated nanowire quantum dots emitting at $\lambda \sim 980$ nm. However, the InAs/InP material system is also the ideal choice for generating telecom wavelengths and we have recently demonstrated high-quality

sources emitting in the O-band [54]. Using such emitters, the pick and place approach described here can be applied to the highly developed silicon-on-insulator integrated photonics platform.

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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On-chip indistinguishable photons using III-V nanowire/SiN hybrid integration: Supplementary Material

Edith Yeung,^{1,2} David B. Northeast,¹ Jeongwan Jin,¹ Patrick Laferrière,¹ Marek
Korkusinski,^{1,2} Philip J. Poole,¹ Robin L. Williams,¹ and Dan Dalacu^{1,2}

¹*National Research Council Canada,
Ottawa, Ontario, Canada, K1A 0R6.*

²*University of Ottawa, Ottawa, Ontario, Canada, K1N 6N5.*

I. THEORY

In this section we derive the relevant formulas used to calculate the coincidences measured in the CW and pulsed Hong-Ou-Mandel experiment. We begin by calculating systematically the second-order correlation functions $g^{(2)}$ for the CW and pulsed excitation schemes. Further, we develop a systematic treatment of the HOM visibility by relating it to the dynamics of the quantum dot by the quantum regression theorem. We then present the formulas $g_{||}^{(2)}(\tau)$ for the HOM visibility in each excitation scheme.

A. Second-order correlation under CW excitation

In the continuous-wave excitation scheme, we assume that we deal with a strictly two-level system excited non-resonantly (incoherently). The higher-energy state, $|1\rangle$, corresponds to the exciton X^{1-} trapped in the quantum dot, whilst the lower-energy state, $|0\rangle$, corresponds to the final-state electron. The state $|1\rangle$ is re-excited at the rate R and recombines radiatively at the rate $\frac{1}{T_1}$, expressed by the radiative lifetime T_1 . We formulate the rate equations for the time-dependent occupation probabilities $p_1(t)$ and $p_0(t)$ of the two levels in the following form:

$$\frac{dp_1(t)}{dt} = -\frac{1}{T_1}p_1(t) + Rp_0(t), \quad (1)$$

$$\frac{dp_0(t)}{dt} = \frac{1}{T_1}p_1(t) - Rp_0(t). \quad (2)$$

The normalization condition $p_1(t) + p_0(t) = 1$ at any time t . By adding the two equations we find that this condition is naturally satisfied. We focus on the first equation and substitute $p_0(t) = 1 - p_1(t)$, obtaining

$$\frac{dp_1(t)}{dt} = -\left(\frac{1}{T_1} + R\right)p_1(t) + R. \quad (3)$$

Postulating the solution in the form

$$p_1(t) = A(t) \exp\left[-\left(\frac{1}{T_1} + R\right)t\right], \quad (4)$$

after elementary calculus we obtain the prefactor $A(t)$ in the form:

$$A(t) = \frac{T_1 R}{1 + T_1 R} \exp\left[\left(\frac{1}{T_1} + R\right)t\right] + A_0, \quad (5)$$

where A_0 is derived from initial conditions. In order to calculate the correlation function $g^{(2)}(t)$, we assume that at $t = 0$ a photon has been emitted, and as a result our system is necessarily in the state $|0\rangle$ and begins to be repopulated. The probability of detecting another photon at time t is proportional to $p_1(t)$ with the above condition. By normalizing it to the probability at $t \rightarrow \infty$ we obtain $g^{(2)}(t)$. We therefore need to calculate the probability $p_1(t)$ with the condition that $p_1(t = 0) = 0$, which determines our choice of A_0 . Elementary calculations give

$$p_1(t) = \frac{T_1 R}{1 + T_1 R} \left\{ 1 - \exp \left[- \left(\frac{1}{T_1} + R \right) t \right] \right\}, \quad (6)$$

$$g^{(2)}(t) = \frac{p_1(t)}{p_1(t \rightarrow \infty)} = 1 - \exp \left[- \left(\frac{1}{T_1} + R \right) t \right]. \quad (7)$$

This formula, generalized to account for negative delays, is used in the main text.

B. Second-order correlation under pulsed excitation

We now develop the formula for the second-order correlation function $g^{(2)}(t)$ under the above-bandgap pulsed excitation scheme. We continue to model the dynamics of the charged exciton in the quantum dot by the two-level system of states $|1\rangle$ (charged exciton present) and $|0\rangle$ (final-state electron), with the radiative recombination rate $1/T_1$ related to the radiative lifetime T_1 . We assume that at time $t = 0$ a strong pulse has been applied nonresonantly, creating a large number $N(t)$ of electron-hole pairs in the sample. The population $N(t)$ is time-dependent since the pairs decay, e.g., by radiative or non-radiative recombination or diffusion. If we characterize the decay process by the time constant T_D , then

$$N(t) = N_0 \exp \left(- \frac{t}{T_D} \right), \quad (8)$$

where N_0 is a constant factor dependent on the power of the excitation pulse (we assume that all electron-hole pairs are created instantaneously). The quantum dot is then refilled from this above-bandgap reservoir at the average rate of $1/\tau_e$ per pair. Therefore, the quantum dot reexcitation rate is

$$R(t) = \frac{N_0}{\tau_e} \exp \left(- \frac{t}{T_D} \right). \quad (9)$$

Here we assume that the capture of one electron-hole pair by the dot alters the population $N(t)$ to a negligible degree. The above time-dependent refill rate can now be inserted into

the rate equations (1) and (2) to obtain the complete time-dependent probability $p_1(t)$ proportional to the photon emission probability. Through the factor N_0 this probability is pulse power dependent. The solution can be obtained in a closed, semi-analytical form, however we find it unwieldy for the analysis in the main text. For any excitation power, we normalize the solution and approximate it with a biexponential formula, which for the n -th pulse takes the form

$$p_1^{(n)}(t) = PH(t - nT) \left[1 - \exp\left(-\frac{t - nT}{\tau_e}\right) \right] \exp\left(-\frac{t - nT}{T_1}\right), \quad (10)$$

with the pulses separated by the fixed time interval T , and $H(t)$ being the Heaviside step function. The factor P is calibrated so that there is only one photon per each pulse, according to the assumptions explained in the main text. We see that, within each pulse, the short-time behavior of this emission probability is described by the reexcitation rate $1/\tau_e$, while the long-time behavior is determined by the radiative lifetime T_1 . However, in this simple phenomenological approach we lose the dependence on excitation power. We can now express the probability of photon emission as a function of time for a train of pulses simply as a superposition

$$p_1(t) = \sum_n p_1^{(n)}(t). \quad (11)$$

The *instantaneous* second-order unnormalized correlation function $g^{(2)}(t_1, t_2)$ expresses the probability of detecting a photon at time t_2 on condition that a photon has been detected at time t_1 . Let us denote the delay $\tau = t_2 - t_1$ and assume that the photon at time t_1 originated from the pulse with index $n = 0$. With these assumptions, and taking into account that only one photon can be emitted as a result of each pulse, we have

$$g^{(2)}(t_1, \tau) = p_1^{(0)}(t_1) \sum_{n \neq 0} p_1^{(n)}(t_1 + \tau). \quad (12)$$

The statistical analysis of the experimental data involves performing an ensemble average of the above function. We can calculate this ensemble average by integrating

$$g^{(2)}(\tau) = \frac{1}{T} \int_{-\infty}^{\infty} g^{(2)}(t_1, \tau) dt_1. \quad (13)$$

With our biexponential probability functions we arrive at an analytical formula

$$\begin{aligned} g^{(2)}(\tau) = & \left(\frac{T_1}{2T} - \frac{1}{T} \frac{\tau_e T_1}{2\tau_e + T_1} \right) \sum_{n \neq 0} \exp\left(-\frac{|\tau - nT|}{T_1}\right) \\ & + \left(\frac{1}{2T} \frac{\tau_e T_1}{\tau_e + T_1} - \frac{1}{T} \frac{\tau_e T_1}{2\tau_e + T_1} \right) \sum_{n \neq 0} \exp\left(-\frac{\tau_e + T_1}{\tau_e T_1} |\tau - nT|\right). \end{aligned} \quad (14)$$

The result (up to a constant factor) is identical to the one obtained by taking the self-convolution of the probability function and removing the part arising from the central peak to account for the single photon per peak emission. This latter procedure is mentioned in the main text.

C. A general theoretical approach to two-photon HOM interference visibility

In this Section we present the general approach to calculating the two-photon HOM interference visibility based on the quantum regression theorem. We follow the method described by Kiraz et al. in Ref. 1 with modifications to the Bloch equations stemming from the specifics of our system.

We begin with the schematic diagram of the experimental system presented in Fig. 3(a) of the main text. If we denote by indices 3 and 4 the output ports of the beamsplitter BS2, then the (unnormalized) HOM cross-correlation function is given by:

$$g_{||}^{(2)}(t, \tau) = \langle a_3^+(t) a_4^+(t + \tau) a_4(t + \tau) a_3(t) \rangle, \quad (15)$$

where $a_m^+(t)$ ($a_m(t)$) is the creation (annihilation) operator of a photon at time t in the mode (channel) m , and the triangular brackets denote the expectation value. To make the HOM interference effects clear, one now expresses the photon creation and annihilation operators at the output ports 3, 4 in terms of those at the input ports 1, 2 of the beamsplitter BS2. For simplicity of presentation we assume BS2 to be a perfect 50% – 50% device; the approach can be readily generalized to any transmission and reflection coefficients, as shown in the main text. We have therefore

$$g_{||}^{(2)}(t, \tau) = \frac{1}{2} \langle a_1^+(t) a_2^+(t + \tau) a_2(t + \tau) a_1(t) \rangle - \frac{1}{2} \langle a_1^+(t) a_2^+(t + \tau) a_1(t + \tau) a_2(t) \rangle, \quad (16)$$

with the crucial HOM interference term appearing with the negative sign. Next, we express the photon operators via the quantum dot raising and lowering (dipole) operators π^+ and π [Ref. 2]. The operators work on the quantum dot states in the following manner: $\pi^+ |0\rangle = |1\rangle$ and $\pi |1\rangle = |0\rangle$, with the state $|1\rangle$ ($|0\rangle$) corresponding to the confined charged exciton, or "filled dot" (the final state electron, or "empty dot"). The equivalence between the photon and quantum dot operators is expressed as^{1,2}

$$a(t) = A(\vec{r}) \pi \left(t - \frac{\vec{r}}{c} \right), \quad (17)$$

where $A(\vec{r})$ is the time-independent proportionality factor, c is the speed of light, and $\vec{r}$ represents the distance (equiv. time delay) between the quantum dot and the beamsplitter. Upon substitution, and omitting the irrelevant time-independent factors, we obtain:

$$g_{||}^{(2)}(t, \tau) = \frac{1}{2} \langle \pi^+(t) \pi^+(t + \delta\tau_p + \tau) \pi(t + \delta\tau_p + \tau) \pi(t) \rangle - \frac{1}{2} \langle \pi^+(t) \pi^+(t + \delta\tau_p + \tau) \pi(t + \tau) \pi(t + \delta\tau_p) \rangle. \quad (18)$$

We note that the two photons are emitted by the same quantum dot, and therefore we only use one set of quantum dot raising and lowering operators, although expressed at different times. The two photons are, in general, not independent, and the ordering of operators matters. We have also accounted for the fact that one of the channels in the setup is guided by the delay line, adding the time shift $\delta\tau_p$. In this manner we have translated the photonic correlation function into a formula expressed by the dynamical properties of the quantum dot. To develop the cross-correlation function further, we therefore have to consider the dynamics of the charged exciton.

The Hamiltonian of the charged exciton confined in the quantum dot, in the language of the dipole operators, is simply

$$\hat{H} = E_1 \pi^+ \pi + E_0 \pi \pi^+, \quad (19)$$

where E_1 is the energy of the charged exciton, and E_0 is the energy of the final-state electron. The system is not driven coherently; the radiative decay process is also incoherent, and therefore is not accounted for in the Hamiltonian. Henceforth we operate in the interaction picture defined with the above Hamiltonian.

The dynamics of our two-level system of charged exciton is described by the master equation

$$\frac{d}{dt} \hat{\varrho} = R[\pi^+ \hat{\varrho}, \pi] + \sum_{n=1}^2 \left(c_n \hat{\varrho} c_n^\dagger - \frac{1}{2} \{c_n^\dagger c_n, \hat{\varrho}\} \right), \quad (20)$$

where $\hat{\varrho}$ is the density matrix, the square (curly) brackets denote the commutator (anticommutator), the collapse operators $c_1 = \frac{1}{\sqrt{T_1}} \pi$ and $c_2 = \frac{1}{\sqrt{2\tau_d}} \sigma_z$, and σ_z is the z Pauli matrix. The master equation is parametrized by the incoherent reexcitation rate R (perhaps time-dependent), the radiative lifetime T_1 and the pure dephasing time τ_d . Let us unfold the master equation into the Bloch equations for the matrix elements of the elements of the

density matrix:

$$\frac{d}{dt}\varrho_{11} = R\varrho_{22} - \frac{1}{T_1}\varrho_{11}, \quad (21)$$

$$\frac{d}{dt}\varrho_{22} = -R\varrho_{22} + \frac{1}{T_1}\varrho_{11}, \quad (22)$$

$$\frac{d}{dt}\varrho_{12} = -\left(\frac{1}{2T_1} + \frac{1}{\tau_d}\right)\varrho_{12}, \quad (23)$$

and the equation for ϱ_{21} is a complex conjugate of that for ϱ_{12} . We identify $\varrho_{11}(t) = p_1(t)$, the probability of the dot being in the "filled" state (charged exciton confined), and $\varrho_{22} = p_0(t)$, the probability of the dot being "empty" (holding just the final-state electron). The time evolution of the coherence ϱ_{12} is parametrized by the coherence time τ_c (also labeled as T_2) defined in the usual manner:

$$\frac{1}{\tau_c} = \frac{1}{2T_1} + \frac{1}{\tau_d}. \quad (24)$$

We note that in the timing of the Bloch equations defined above, our approach differs somewhat from that presented in Ref. 1. The equations (21), (22) are coupled and have to be solved as a set; we have already shown these solutions for the CW and pulsed excitation cases in the previous Sections. The equation (23) for the coherence will now be used to develop the interference term of the HOM two-photon visibility.

We now return to the HOM visibility, Eq. (18) and focus on the first term. It expresses the probability of emission of the photon at time $t + \delta\tau_p + \tau$ on condition that a photon has been emitted at time t . We track the pair of dipole operators dependent on the time τ , related to the time in which the two photons interfere on the beamsplitter. We use the quantum regression theorem^{1,2}. Since the expectation value of the product $\pi^+(t+t')\pi(t+t')$ is the probability $p_1(t+t') = \varrho_{11}(t+t')$ (we denote $t' = \delta\tau_p + \tau$), we need to solve the Bloch equations (21) and (22) with the initial condition

$$\langle \pi^+(t)\pi^+(t+t')\pi(t+t')\pi(t) \rangle|_{t'=0} = \langle \pi^+(t)\pi^+(t)\pi(t)\pi(t) \rangle. \quad (25)$$

In the case of independent single-photon sources (e.g., two independent quantum dots) the above value can be nonzero. However, we implement our HOM experiment with one single-photon source, for which the above probability of emitting two photons at the same time t is zero. As a result of this recipe we simply obtain the instantaneous second-order correlation function $g^{(2)}(t, \tau)$ of the charged exciton in the quantum dot (up to normalization). In the

previous Sections we have derived the analytical formulas for $g^{(2)}(t, \tau)$ both in the CW and pulsed excitation.

Next, we consider the interference term $\langle \pi^+(t)\pi^+(t + \delta\tau_p + \tau)\pi(t + \tau)\pi(t + \delta\tau_p) \rangle$. Here we see four dipole operators, each at a different time, however only two operators depend on the crucial time delay τ . Further, since each of them is taken at a different time (with the time difference $\delta\tau_p$), we have to treat these operators separately. The expectation value of the operator π^+ (π) is related to the coherence ϱ_{21} (ϱ_{12}) and therefore the equation of motion for the interference term is

$$\frac{d}{d\tau} \langle \pi^+(t)\pi^+(t + \delta\tau_p + \tau)\pi(t + \tau)\pi(t + \delta\tau_p) \rangle = -\frac{2}{\tau_c} \langle \pi^+(t)\pi^+(t + \delta\tau_p + \tau)\pi(t + \tau)\pi(t + \delta\tau_p) \rangle. \quad (26)$$

The initial condition for the interference term is evidently (up to normalization factor)

$$\begin{aligned} \langle \pi^+(t)\pi^+(t + \delta\tau_p + \tau)\pi(t + \tau)\pi(t + \delta\tau_p) \rangle|_{\tau=0} &= \langle \pi^+(t)\pi^+(t + \delta\tau_p)\pi(t)\pi(t + \delta\tau_p) \rangle \\ &= g^{(2)}(t + \delta\tau_p). \end{aligned} \quad (27)$$

After elementary calculus we find

$$\langle \pi^+(t)\pi^+(t + \delta\tau_p + \tau)\pi(t + \tau)\pi(t + \delta\tau_p) \rangle = g^{(2)}(t + \tau_p, \tau = 0) \exp\left(-\frac{2}{\tau_c}\tau\right). \quad (28)$$

We have arrived at the final expression for the instantaneous interference visibility:

$$g_{||}^{(2)}(t, \tau) = g^{(2)}(t, \delta\tau_p + \tau) - g^{(2)}(t, \delta\tau_p) \exp\left(-\frac{2}{\tau_c}\tau\right). \quad (29)$$

We note that the instantaneous second-order correlation function $g^{(2)}$ appears in both terms, although it is taken at a different time. This allows us to normalize both terms with an identical factor. Also, at the delay $\tau = 0$, $g_{||}^{(2)}(t, \delta\tau_p) = 0$ identically, which agrees with the previously published results in the electromagnetic domain³⁻⁶. The closed formulas for $g_{||}^{(2)}(t, \tau)$ and its ensemble average depend on the excitation conditions and we will summarize them below.

D. Two-photon interference visibility under CW excitation

Let us illustrate the use of the formula (29) on the example of a system with the CW excitation, beginning with a system of two independent single-photon sources with identical coherence properties¹. In this case, the (unnormalized) conditional probability of photon

emission is nonzero and independent of time, i.e., $g^{(2)}(t, \tau) = p_1 p_2$, where p_1 is the probability of emission from the first source, and p_2 is the probability of emission from the second source. In such case, the normalized HOM visibility attains the well-known form

$$g_{||}^{(2)}(t, \tau) = g_{||}^{(2)}(\tau) = \frac{1}{2} \left[1 - \exp \left(-\frac{2}{\tau_c} \tau \right) \right]. \quad (30)$$

Let us now consider the system described in the main text, consisting of one single-photon source, with one of the photons guided through the mode imposing the delay $\delta\tau_p$. We assume further that the system has reached its steady state, i.e., the second-order correlation function does not depend on the time t of emission of the first photon of the pair, but depends only on the overall delay $\delta\tau_p + \tau$. If the delay $\delta\tau_p$ is sufficiently long, then we may approximate

$$g^{(2)}(\delta\tau_p) \approx g^{(2)}(\delta\tau_p + \tau) \quad (31)$$

and the HOM visibility becomes

$$g_{||}^{(2)}(t, \tau) = g_{||}^{(2)}(\tau) = \frac{1}{2} g^{(2)}(\delta\tau_p + \tau) \left[1 - \exp \left(-\frac{2}{\tau_c} \tau \right) \right]. \quad (32)$$

The HOM visibility has a factor originating from the single-photon nature of the photon source. However, owing to the two emission events being separated by a long delay (at least $\delta\tau_p$), the photons are approximately uncorrelated, and the system has coherence characteristics as if it involved two independent photon sources. This assumption is made in the main text, however, it is adjusted to account for the subtleties of the two-beamsplitter setup.

Finally, if the delay $\delta\tau_p$ is not long (comparable to the coherence time τ_c), the complexities of conditional two-photon emission have to be taken into account, and we have

$$g_{||}^{(2)}(t, \tau) = g_{||}^{(2)}(\tau) = \frac{1}{2} \left[g^{(2)}(\delta\tau_p + \tau) - g^{(2)}(\delta\tau_p) \exp \left(-\frac{2}{\tau_c} \tau \right) \right]. \quad (33)$$

In the two latter cases we utilize the formula (7) for $g^{(2)}(\tau)$ obtained in Section I A.

E. Two-photon interference visibility under pulsed excitation

Under pulsed excitation, the reexcitation rate R in Bloch equations (21), (22) is time-dependent and periodic. In our system, the delay $\delta\tau_p$ between the two photons is chosen to correspond precisely to the time interval T between pulses. As a result, we first consider

the instantaneous HOM visibility expressed by Eq. (29), with the instantaneous second-order correlation function $g^{(2)}(t_1, \tau)$ given by Eq. (12). To enable comparison with the experimental data, we perform the ensemble average:

$$g_{||}^{(2)}(\tau) = \frac{1}{T} \int_{-\infty}^{\infty} g_{||}^{(2)}(t_1, \tau) dt_1. \quad (34)$$

We obtain

$$g_{||}^{(2)}(\tau) = \frac{1}{2} \left[g^{(2)}(T + \tau) - g^{(2)}(T) \exp\left(-\frac{2}{\tau_c} \tau\right) \right], \quad (35)$$

i.e., the formula equivalent to that used in CW excitation with short delays $\delta\tau_d$. However, here the second-order correlation function $g^{(2)}(\tau)$ attains the far more complex, time-periodic form (14) developed in Sec. IB. This formula, modified to account for the subtleties of the two-beamsplitter setup, is used in the main text.

We conclude this presentation by noting that the case of pulsed excitation of two independent single-photon sources was considered by Kiraz et.al.¹. The model described therein required a much simpler calculation owing to the possibility of separating the operator time evolution into two pairs (for each emitter separately). Nevertheless, the results obtained are similar to those reported here. This is due to the fact that in our system, each pulse resets the dynamics of the quantum dot. As a result, the HOM interference of two photons emitted in our quantum dot after different pulses is formally equivalent to that of two photons produced by independent, but identical emitters. Our system is nontrivially different only if both photons can be emitted after the same pulse (although not at the same time). This case is, however, beyond the scope of this work.

II. ADDITIONAL FIGURES

In the following we provide additional figures supporting the experiments in the main text.

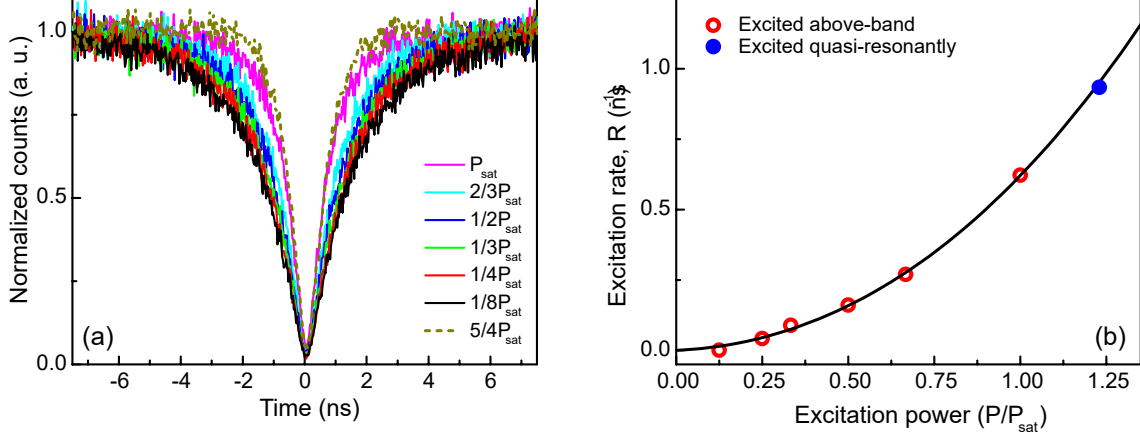


FIG. 1: (a) Second-order correlation, $g^{(2)}(\tau)$, as a function of CW excitation power. Solid (dashed) lines correspond to measurements made using above-band (quasi-resonant) excitation. (b) Extracted excitation rate, R , using $g^{(2)}(\tau) = 1 - e^{-(1/T_1+R)|\tau|}$ where T_1 is the radiative lifetime. The formula is derived in Section I.

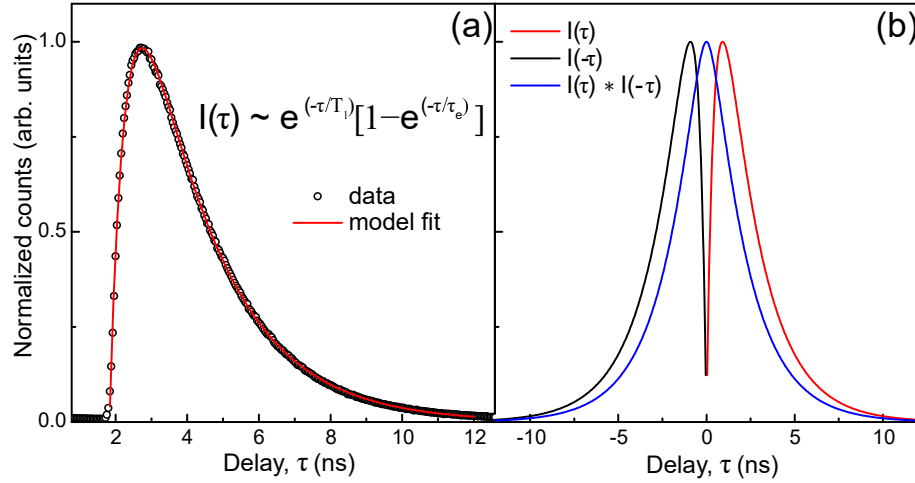


FIG. 2: (a) Time-resolved photoluminescence of the X^{1-} emission excited quasi-resonantly at $P = P_{\text{sat}}$ (open symbols). The radiative lifetime, T_1 , and the state preparation time, τ_e , is extracted from the model fit of the PL decay (red curve). (b) Temporal profile of a $g^{(2)}(\tau)$ side peak constructed from the self-convolution of the model fit to the radiative decay. This approach is equivalent to the ensemble-averaged calculation of conditional probabilities for the two photons described in Section I.

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4.2 Conclusion

This chapter laid the groundwork for the type of experiments performed in subsequent chapters. The methods and results from this chapter demonstrates a robust pathway for on-chip generation of indistinguishable single-photons which is a key requirement in many quantum photonic applications. The hybrid integration of III–V nanowire quantum dots with Si_3N_4 waveguides bridges two previously distinct technologies: high-quality single-photon emitters and mature silicon-based photonic platforms.

The results are significant for the broader field of quantum photonic technologies because they show that solid-state quantum emitters can be reliably integrated into planar photonic integrated circuits without compromising their optical qualities. By achieving high coupling efficiencies, bright emission, and measurable photon indistinguishability directly on-chip, this platform provides a scalable route for building more complex photonic circuits. This will be demonstrated later in Ch. 5 where the NW-QDs are integrated with more complex photonic circuitry for widening applications.

Chapter 5

Advanced Architectures On-chip

In the previous chapter, we have demonstrated that the hybrid-devices can be good sources of single-photons using a simple straight waveguide and a sufficiently tapered InP nanowire. In this chapter, we expand on this by incorporating different architectures on-chip. As mentioned in the previous chapter, the collection efficiency from the tapered end of the nanowire is at least ten times higher than from the untapered end of the nanowire. In order to recuperate the other half of the photons, we employ a different architecture that will allow the QD emission to efficiently couple into the waveguide from both ends of the nanowire.

Using the same principles, we evanescently couple the QD emission using an untapered nanowire and a curved waveguide. The goal is to recuperate the lost photons emitting in the opposite direction by collecting QD emission from both ends of the device.

5.1 Article: Advanced architectures for coupling III-V nanowires to photonic integrated circuitry

Adapted with permission from **Yeung, E.**, Sorensen, K., Northeast, D. B., Milanizadeh, M., Poole, P. J., Williams, R. L., Dalacu, D. (2026.) Nanowire Quantum Dot Emitters Evanescently Coupled On-Chip via Curved and Segmented SiN Waveguides. *ACS Photonics*, 13(15), 4185-4192. <https://doi.org/10.1021/acsp Photonics.6c00397>. Copyright 2026 American Chemical Society (ACS).

5.1.1 Contributions of co-authors

Design and simulations for the PIC chips was done by David B. Northeast and Maziyar Milanizadeh. The design was sent off to a commercial foundry, LIGENTEC, for fabrication. The nanowires were grown by Philip J. Poole.

Both the nanowires and chip was pre-characterized by Edith Yeung prior to transfer onto the PIC chip. E. Yeung also picked the nanowire from its growth substrate and transferred it to the devices presented in this chapter.

Experimental design, setup, measurements, and analysis were performed by E. Yeung with the assistance of K. Sorensen. D. Northeast supervised the experimental aspects of the project. Both claim equal contribution for the experimental work where E. Yeung was mentoring and the co-op student, K. Sorensen, through the experimental process and handling of laboratory equipment. Required fibre lengths for the coherence measurement setup was calculated by E. Yeung. J. F. Weber helped splice the PM fibres.

Bi-directional experiments were conceived by E. Yeung. Code to run the coherence measurements was written by E. Yeung.

E. Yeung wrote the manuscript. D. Dalacu, P. J. Poole, and R. L. Williams conceptualized the original project. All authors discussed the results and contributed to the text of the manuscript.

D. Dalacu supervised the research.

Advanced architectures for coupling III-V nanowires to photonic integrated circuitry

Edith Yeung,^{1,2} Kataryna Sorensen,^{1,2} David B. Northeast,² Maziyar Milanizadeh,² Philip J. Poole,² Robin L. Williams,² and Dan Dalacu^{1,2}

¹*University of Ottawa, Ottawa, Ontario, Canada, K1N 6N5.*

²*National Research Council Canada, Ottawa, Ontario, Canada, K1A 0R6.*

This work implements a hybrid device based on a semiconductor quantum dot embedded within a nanowire to bridge a non-continuous curved waveguide structure. The geometry takes advantage of evanescent coupling between the photonic structures to recover single photons emitted from both outputs of the device. Auto- and cross-correlation measurements were performed on different output facets of the device. We demonstrate single-photon emission from both ends of the nanowire for both neutral, X and XX , and charged X^- , excitonic complexes. We further demonstrate the cascaded $XX - X$ emission by collecting each complex from a different facet. This work lays the foundation for on-chip architectures which utilize multi-directional integration of quantum emitters.

I. INTRODUCTION

Sources of single-photons are a key resource in quantum information processing applications¹ and solid-state two-level systems² offer a promising route to generating single-photons on-demand with high efficiency. Semiconductor quantum dots (QDs) are of particular relevance as they have achieved nearly ideal performance when integrated with photonic structures^{3–9}. These structures, however, are typically designed to operate out-of-plane whereas future applications will likely require the scalability provided by on-chip integration^{10–13}.

QDs embedded within nanowires that are grown¹⁴ (i.e. bottom-up) as opposed to etched¹⁵ (i.e. top-down) provide a route to fabricating sources of single photons deterministically. The number of emitters per device is controlled and each device is grown at a specified position on the substrate. Bottom-up nanowire QDs have been used to generate both high quality single photons^{16,17} and high fidelity entangled photon pairs^{18,19}. These sources were also operated out-of-plane and high efficiency was obtained by designing the nanowire to act as a single mode waveguide supporting only the fundamental HE_{11} mode^{20,21} and by introducing a taper along the length of the nanowire²² to expand the mode for efficient coupling to external optics. Properly designed, 95% of the dot emission is directed into the HE_{11} mode²¹ which can be coupled to single-mode fibre with an efficiency of 93%²³.

The ability to control the nanowire taper is also relevant for designing chip-integrated sources^{24,25}. In this case, the mode expansion provided by the taper is not for efficient coupling to low numerical aperture optics, but rather for efficient evanescent coupling to on-chip waveguides. By placing an appropriately tapered nanowire along a straight Si_3N_4 waveguide, as in Fig. 1(a), the HE_{11} mode can be efficiently coupled on-chip, with simulated coupling efficiencies approaching 100%^{24,25}. The hybrid structures can be assembled by picking up nanowires from the growth substrate and placing them on Si_3N_4 waveguide using a scanning electron microscope (SEM)-based nanomanipulator^{24,26}. With this approach, we have demonstrated >90% coupling of the HE_{11} mode

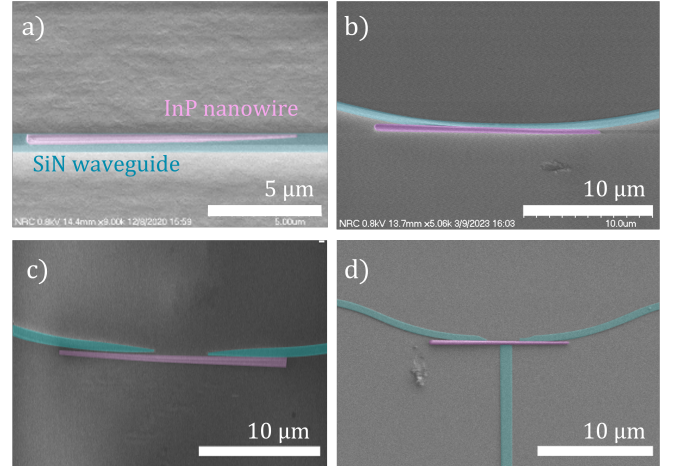


FIG. 1. False colour SEM images with SiO_2 in grey, Si_3N_4 in blue and InP in purple. (a) Straight Si_3N_4 waveguide with a tapered InP nanowire on top. (b) Curved waveguide with a nanowire on the side. (c) Segmented waveguide with a nanowire bridging the gap. (d) As in (c) but with an additional waveguide for cross-pumping.

to a Si_3N_4 waveguide and have generated on-chip single-photons with purities exceeding 95%²⁵.

The pick and place hybrid platform is not limited to simple structures such as that shown in Fig. 1(a): one can also construct hybrid devices consisting of a nominally untapered nanowire with a thinner diameter placed beside a curved waveguide as shown in Fig. 1(b). In this case the quantum dot within the nanowire is located at its midpoint and the evanescent transfer of the nanowire mode to the waveguide occurs as the separation between the two increases. Such a device can harvest photons emitted in either direction along the nanowire waveguide, a limitation of the unidirectional structure in Fig. 1(a) where emission towards the base is not efficiently collected.

Expanding on the curved waveguide architecture, a gap can be introduced in the waveguide with the nanowire connecting the circuit, as shown in Fig. 1(c). In this case, the waveguide mode can be transferred completely to the

nanowire and back again, assisted by tapering each end of the segmented waveguide. This structure is relevant for new types of experiments such as transmission measurements to investigate coherent interactions between single photons and two-level systems²⁷. By incorporating an additional waveguide orthogonal to the nanowire, as in Fig. 1(d), we can also envision a resonant fluorescent excitation scheme where pump laser rejection is achieved using orthogonal excitation and collection directions^{28–30}.

In this article we explore some of the advanced architectures shown in Fig. 1, specifically (b) and (c), for on-chip single-photon generation in Si-based integrated photonic circuitry i.e. Si_3N_4 waveguides. We evaluate, through simulations, geometry-dependent coupling efficiencies between a nanowire and a curved waveguide, both with and without a gap. We find the segmented waveguide design more robust against variations in the nanowire diameter when collecting from curved waveguide geometries and use this architecture to demonstrate efficient generation of on-chip high purity single photons. We further utilize the bi-directional nature of the device to perform cross-correlation measurements by collecting from both ends of the nanowire with the nanowire acting as an integrated beamsplitter.

II. SIMULATIONS

As mentioned in the introduction, a nanowire containing a quantum emitter can transfer its light efficiently to a nearby waveguide by evanescent coupling^{24,25}. This process is robust provided both structures allow for a single mode and a tapered section is included over a sufficient length to allow adiabatic transfer of mode energy between the two. In this scenario, light that is emitted away from the taper—50% of light from the quantum dot—will encounter an abrupt end of the nanowire and will couple poorly into the waveguide. A waveguide that curves as it approaches the nanowire will be seen as a gradual change in the refractive index. Device designs can use this effect to allow for light coupling in both directions. Since quantum dot emission has a narrow bandwidth and optical response of all structures is expected to change slowly with wavelength, frequency domain finite element simulations (COMSOL Multiphysics) are performed to predict device performance.

We first investigate power transfer from a dipole in a straight InP nanowire to a curved Si_3N_4 waveguide. A $20\text{ }\mu\text{m}$ nanowire length is used in order to provide sufficient separation between the nanowire ends and the curved waveguide for a waveguide with a $30\text{ }\mu\text{m}$ radius of curvature. The Si_3N_4 waveguides are $0.4\text{ }\mu\text{m}$ thick and $0.5\text{ }\mu\text{m}$ wide, chosen to support a single mode in both TE and TM polarizations, sitting on $4\text{ }\mu\text{m}$ of SiO_2 . The nanowire, with radius r_{nw} , is placed as seen in Fig. 2(a) such that the dipole is located precisely where the nanowire meets the waveguide. The figure shows the electric field polarized along the TE direction of the

waveguide. In the inset, we show the fraction of the total power radiated from the dipole that couples to the waveguide for different r_{nw} . We calculate a peak coupling of greater than 76% in both waveguide directions which drops quickly for nanowire radii away for the optimal value of $r_{\text{nw}} = 115\text{ nm}$.

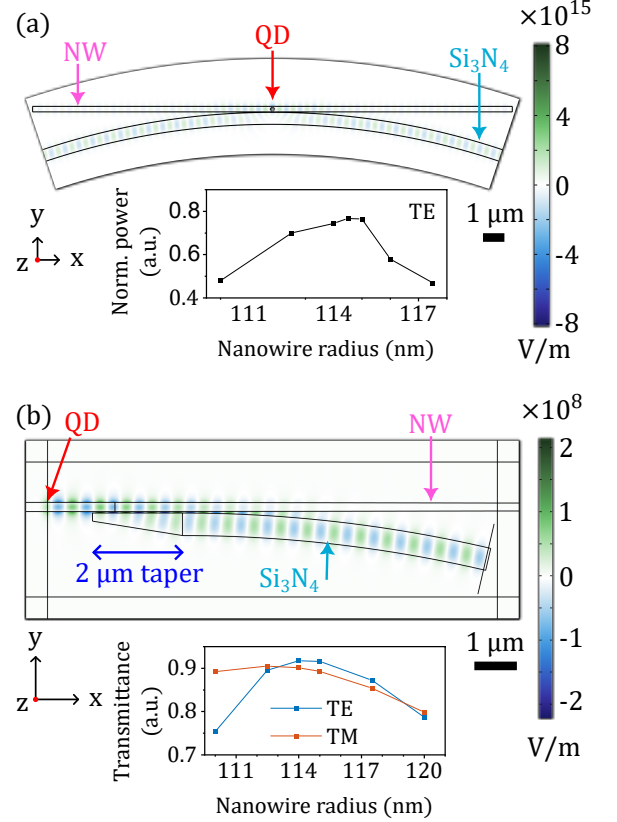


FIG. 2. Finite element simulations were performed for the two different hybrid coupling approaches. (a) Top-down view of the photonic integrated circuit. The quantum dot emitter (here a dipole) is located at the midpoint of the InP nanowire (NW) which is tangent to the curved Si_3N_4 waveguide. The emission from the quantum dot is emitted into the nanowire and couples into the waveguide as it bends away. The total power coupled into the waveguide as a function of nanowire radius, r_{nw} , is shown in the inset. (b) Revised coupling design that makes use of a gap in the Si_3N_4 waveguide with tapered sections at the gap. The inset shows the improved transmittance between both HE_{11} polarizations in the nanowire and the TE/TM modes in the Si_3N_4 ridge waveguide. Electric field simulations in both (a) and (b) are shown for optimal nanowire radii.

Combining the curved geometry with a tapered waveguide, power transfer can be made less sensitive to the nanowire radius, r_{nw} . Fig. 2(b) depicts a curved Si_3N_4 waveguide that tapers to an end over a length of $2\text{ }\mu\text{m}$. Simulated transmittance of HE_{11} nanowire modes to TE and TM waveguide modes is shown in the inset for different r_{nw} . We calculate a peak transmission of over 90% for both linear polarizations and observe a reduced de-

pendence on r_{nw} compared to the untapered waveguide geometry in Fig. 2(a).

In summary, these simulations show that curved waveguide geometries allow for the efficient collection of quantum dot emission directed towards either end of the InP nanowire. This is in contrast to tapered nanowires on straight waveguides where coupling efficiency for emission towards the base of the nanowire is low and scattering at the base is high^{24,25}. When comparing the continuous curved waveguide with the segmented waveguide, the simulations show that segmented tapered Si_3N_4 geometry provides higher coupling values with a higher tolerance for variations in nanowire radius. In the experiments below, we will focus on the segmented tapered Si_3N_4 geometry.

III. EXPERIMENTAL DETAILS

The quantum emitters used in this work are InAsP quantum dots embedded within InP nanowires. An array of position-controlled InP nanowires are grown using selective-area vapour-liquid-solid epitaxy, see Refs. 14,16, 17 for details. In the first growth stage, we grow a single InAsP QD as a segment within a 20 nm diameter InP nanowire core at a nominal height of 10 μm above the substrate. In the second growth stage, the nanowire core is clad with InP to increase the diameter to 210 nm and to increase the nanowire height to 20 μm such that the QD is at the midpoint of the nanowire. In this cladding stage, growth conditions are adjusted to minimize tapering of the nanowire diameter along its length as the proposed designed is based on an untapered nanowire, see Fig. 2.

The integrated photonic circuitry is based on low-pressure chemical vapour deposited Si_3N_4 waveguides fabricated on Si wafers using the commercial foundry Ligentec. Waveguide dimensions are 0.4 μm tall and 0.5 μm wide clad with 3.3 μm of SiO_2 above and 4 μm below. For coupling on and off chip, we use edge couplers designed for lensed fibres: at the etched facets, the waveguide widths are tapered to 0.2 μm over a length of 150 μm . The segmented waveguide consists of a curved section with a 30 μm radius of curvature and with a 4 μm gap at the midpoint. On either side of the gap, the waveguide is tapered down to 0.180 μm over 2.0 μm , as shown in Fig. 2(b). A window in the top SiO_2 layer is opened to expose just the segmented waveguide section, permitting nanowire placement. The etch depth of the SiO_2 exposing part of the Si_3N_4 waveguide was chosen to allow the nanowire to sit roughly at the midpoint (in height) of the waveguide. We note that the asymmetry in the z-direction within the window defines the two orthogonally polarized HE_{11} nanowire modes: one along z and the other parallel to the substrate, i.e. corresponding to the TE and TM modes, respectively, of the Si_3N_4 waveguide.

The pre-characterized nanowire is picked up from the growth substrate with a piezo-motor controlled tungsten

tip on the nanomanipulator inside an SEM^{24,25}. The nanowire is placed tangential to the curved waveguide such that it bridges the gap with the midpoint of the nanowire, and hence the quantum dot, centred in the gap. An example of the final structure, which we term a gap-coupled device, is shown in Fig. 1(c).

The device was measured in a fibre-coupled closed-cycle helium cryostat at 5.5 K. The cryostat is equipped with three xyz piezo stacks: one for free-space alignment and two for aligning fibres to edge facet couplers. We take advantage of the on-chip platform to pump the quantum dot and collect the emission in three different configurations: (i) pump via fibre through the waveguide from one side and collect from the other side, (ii) pump and collect from the same side using a 90:10 fibre splitter (here we can also simultaneously collect from the other side), and (iii) pump from the top using free-space optics and collect from both sides of the segmented waveguide. For coupling off- and on-chip, we used PM 780HP lensed fibres with a mode field diameter of $5.3 \pm 0.1 \mu\text{m}$ at 850 nm. For free-space pumping and for imaging, a 100x objective ($\text{NA} = 0.81$) placed inside the cryostat for used.

Photoluminescence (PL) measurements were done using above-band continuous wave (CW) or pulsed laser excitation. The emission was collected using one of the above configurations and sent to a fibre-coupled grating spectrometer for spectrally-resolved measurements or to a fibre-coupled narrow-band (bandwidth $BW = 0.1 \text{ nm}$) tuneable filter for measurements on single excitonic lines. For measurements on the single excitonic complexes, including time-resolved PL, high-resolution PL and time-correlated PL, emitted photons were detected using superconducting nanowire single-photon detectors (SNSPDs) with efficiencies of 80% and timing jitters of 100 ps.

The $g^{(1)}(\tau)$ measurements were performed using a CW 808 nm laser using setup (iii) described above. To balance the two arms of the interferometer, an in-line variable optical attenuator (VOA) was calibrated to balance the count rate between the two arms before the last beam-splitter.

IV. RESULTS

In Fig. 3 we show a PL spectrum from one exemplary gap-coupled device shown in Fig. 1(c). The measurement was performed by exciting through the waveguide from one side and collecting from the other side. The spectrum shows three dominant peaks, typical in this quantum dot system (see, for example, Ref. 31). The peaks are identified as emission from the neutral exciton (X), biexciton (XX) and trion (X^-) where peak identification is based on the intensity dependence on pump power (see inset) as well as high-resolution PL and cross-correlation measurements discussed below.

The integrated intensities from each complex for pulsed excitation at 80 MHz at saturation, P_{sat} , were 575 kcps

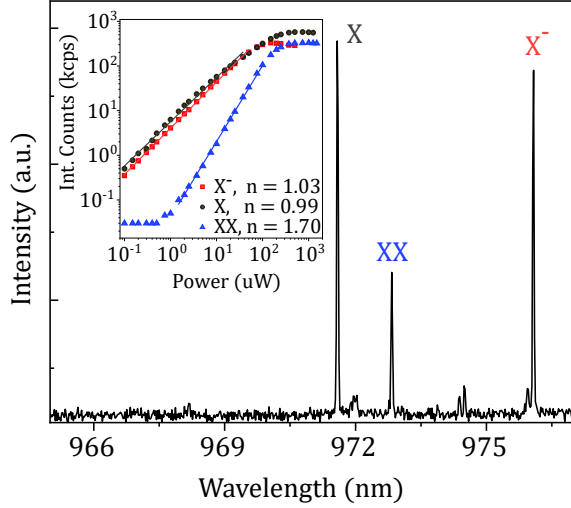


FIG. 3. PL spectrum measured on a gap-coupled device showing three dominant peaks identified as the excitonic complexes X , XX and X^- . The spectrum was obtained with above-band excitation at 780 nm using a pulsed laser operating at an 80 MHz repetition rate with an excitation power of $150 \mu\text{W}$ ($P = 0.7P_{\text{sat}}$). The inset shows the integrated intensity as a function of pulsed excitation power at 80 MHz for each complex: linear ($n \sim 1$) for X (black circles) and X^- (red squares) and approaching quadratic ($n \rightarrow 2$) for XX (blue triangles)³².

for X , 330 kcps for XX and 330 kcps from X^- . Using these values, we obtain end-to-end efficiencies (detected counts/80 MHz) of 0.72%, 0.41%, and 0.41% for X , XX and X^- , respectively. These efficiencies are limited predominantly by the throughput of the setup ($< 10\%$) and by the loss of 50% of the photons directed towards the pump direction as here we are only collected from one side. Additional losses include the nanowire to waveguide coupling, waveguide propagation loss, facet coupling and the detector efficiency, see Supplemental Materials for a detailed loss analysis.

If we reverse the excitation and collection waveguides, we observe a decrease in the count rate of two orders of magnitude whereas simulations predict equal coupling in both directions. This asymmetry may be due to a displacement of the QD from the centre of the gap in the waveguide, either from the dot not centered in the InP nanowire and/or the nanowire not centred in the gap, see supplementary material. A direction-dependent coupling may also arise from an unintentional taper of the nanowire that can be observed in Fig. 1(c). Nonetheless, we obtain sufficient counts from both directions to perform the bi-directional measurements that will be discussed below.

Time-resolved PL traces of the three complexes are shown in Fig. 4(b,d,f) where we excite from one side and collect for the other as above, see Fig. 4(a). To extract decay rates, we assume a double exponential of the form³³

$$I(\tau) = A_b e^{-\frac{\tau}{\tau_b}} + A_d e^{-\frac{\tau}{\tau_d}} + y_0 \quad (1)$$

where τ is the delay between a synchronization pulse from the laser and a detection event, $1/\tau_b$ is the radiative decay rate of the bright exciton, $1/\tau_d$ is the spin-flip rate that converts the system from a dark to bright state with A_b and A_d the amplitudes of the bright and dark contributions, respectively, and y_0 is a background offset. The ratio A_b/A_d is a measure of the occupation probability of the bright state. Model fits to the decay curves are shown as red lines in Fig. 4(b,d,f) and extracted parameters τ_b , τ_d and A_b/A_d are summarized in Table. I for all three complexes.

The time-resolved PL traces of the X and X^- emission, Fig. 4(b) and (f), respectively, both show strong slow decay contributions with decay times, τ_d , consistent with spin-flip rates of neutral excitons in quantum dot systems³³. For X^- , the spin-flip is associated with a carrier in a p-shell of the dot: a dark excited state of the complex that relaxes to the bright state¹⁷. In contrast, the PL decay of the XX emission, Fig. 4(d), is predominantly mono-exponential with a possible small contribution from a slightly slower process.

Second-order autocorrelations, $g^{(2)}(\tau)$, measured using pulsed diode excitation ($\lambda = 670 \text{ nm}$, pulse width = 100 ps) in a Hanbury-Brown and Twist (HBT) setup, see Fig. 4(a), are shown in Fig. 4(c,e,g) for X , XX and X^- , respectively. Due to the significant slow component in the decay of X and X^- , an excitation repetition rate of 10 MHz was used to allow the system to fully decay between pulses. For XX , where τ_d was significantly shorter, the dot was excited at 40 MHz. The reduced zero-delay peak in the $g^{(2)}(\tau)$ for all three complexes demonstrates single photon emission. To quantify the single photon purity, the $g^{(2)}(\tau)$ is modeled by

$$g^{(2)}(\tau) = C_0 \left[e^{-\frac{|\tau|}{\tau_b}} + \frac{A_d}{A_b} e^{-\frac{|\tau|}{\tau_d}} \right] - B_0 e^{-\frac{|\tau|}{\tau_{\text{re}}}} + C_1 \sum_{n \neq 0} \left[e^{-\frac{|\tau| - nt_{\text{rep}}}{\tau_b}} + \frac{A_d}{A_b} e^{-\frac{|\tau| - nt_{\text{rep}}}{\tau_d}} \right] + C_b \quad (2)$$

where C_0 is the amplitude of the coincidences due to the multiphoton emission events at $\tau = 0$, B_0 is associated with re-excitation of the dot, at a rate of τ_{re} , from the same excitation pulse, C_1 corresponds to coincidences recorded from different excitation pulses at $\tau = nt_{\text{rep}}$ for $n \neq 0$, and C_b represents the background counts. For the remaining parameters, values extracted from the fits to the time-resolved PL are used. The model fits are shown in Fig. 4(c, e, g) (red curves) for the three complexes. From the ratio of the area of the zero delay peak to the average peak area at $\tau = nt_{\text{rep}}$, we calculate the probability of multiphoton emission, $g^{(2)}(\tau = 0)$, listed in Table. I.

We obtain very low probabilities of multi-photon emission from X and X^- with $g^{(2)}(\tau = 0)$ values of $< 4\%$ and $< 1\%$, respectively. Coincidences in the zero-delay peak are associated with re-excitation of the dot from the same laser pulse, previously seen in this system with

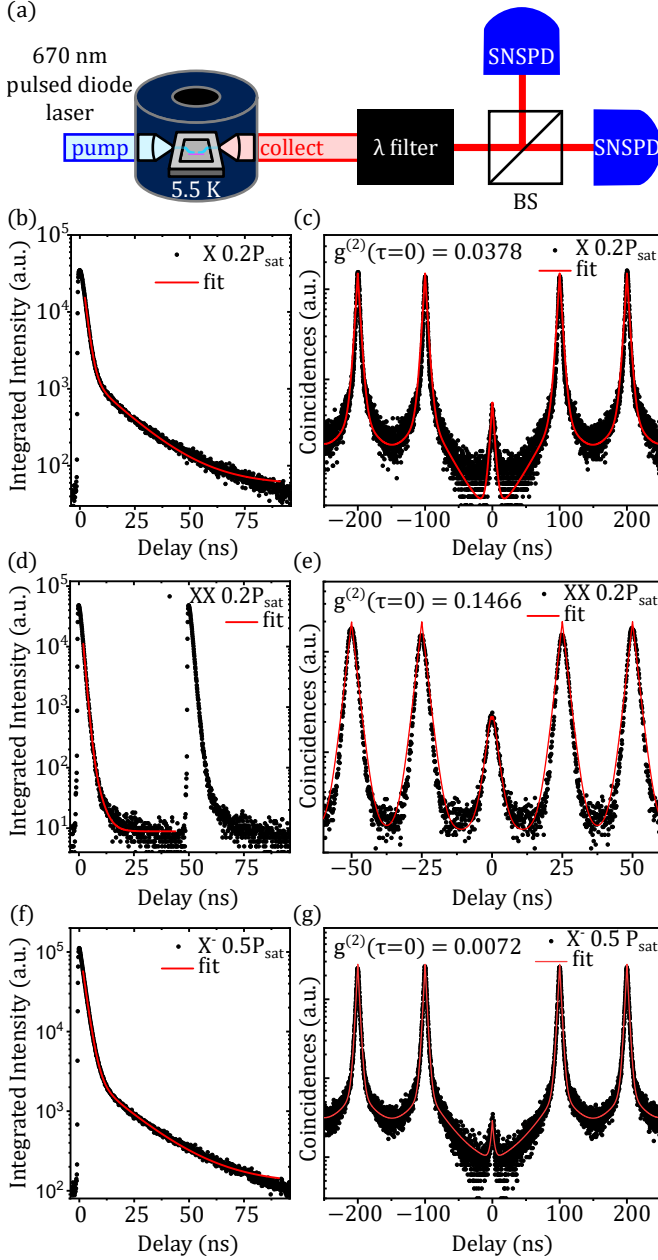


FIG. 4. (a) Schematic of the setup for measuring $g^{(2)}(\tau)$. The excitonic complex to be measured is first filtered then directed to a beamsplitter with the output ports sent to two SNSPD detectors. (b, d, f) PL decay traces and (c, e, g) auto-correlation measurement for the three complexes, X , XX , and X^- , respectively, plotted semi-logarithmically. The excitation rate used in each measurement is shown in the figure as a fraction of the saturation power, P_{sat} .

Complex	τ_b (ns)	τ_d (ns)	A_d/A_b	$g^{(2)}(\tau=0)$
X	1.680 ± 0.004	16.2 ± 0.2	0.0296	0.03776
XX	0.982 ± 0.006	2.8 ± 0.3	0.011	0.14663
X^-	2.107 ± 0.002	18.2 ± 0.2	0.046	0.00724

TABLE I. Summary of parameters extracted from model fits to the PL decay traces and second-order correlation measurements.

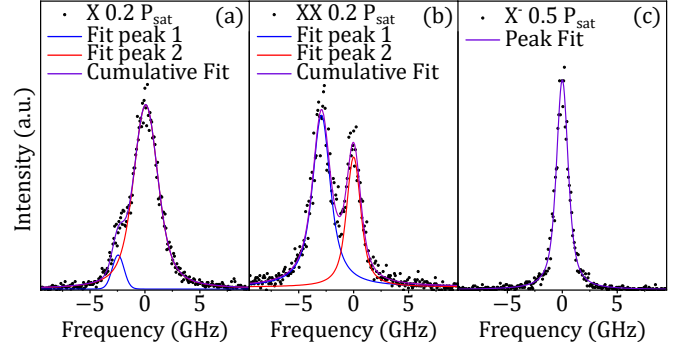


FIG. 5. High-resolution PL of the (a) X , (b) XX and (c) X^- peaks measured at excitation powers indicated. Model fits are shown in purple. For the neutral complexes, the two peaks that constitute the doublets are shown in blue and red.

above-band excitation^{34–36}. We note that the dip in the zero-delay peak that is characteristic of re-excitation is not observed and we speculate that the timescales associated with the re-excitation process are too fast to be resolved using our 100 ps detectors. The higher $g^{(2)}(\tau=0)$ for XX of nearly 15% may be related to the significantly faster decay ($\tau_b = 982$ ps) which would lead to an increased re-excitation probability during the 100 ps excitation pulse.

To evaluate the spectral purity of the emitted photons, we measure high-resolution PL spectra for each complex by directing the filtered emission through a tuneable Fabry-Perot etalon with a bandwidth of $BW = 200$ MHz, sufficiently narrow to resolve real lineshapes. High-resolution spectra of X and XX are shown in Fig. 5(a) and (b), respectively. In each case, two peaks are resolved corresponding to the two orthogonal exciton states separated by the fine structure splitting, FSS ³⁷. The peaks are fit with Voigt lineshapes from which we extract an $FSS = 2.55 \pm 0.02$ GHz for X and 2.96 ± 0.01 GHz for XX . These values of FSS are larger than we typically observe in this system¹⁸ which we associated with slower switching between Group V precursors during the dot growth in this sample that results in an axial asymmetry due to P and As tailing³⁸.

The extracted linewidths are $\delta\omega_{V,X_1} = 1.83 \pm 0.07$ GHz and $\delta\omega_{V,X_2} = 2.59 \pm 0.03$ GHz for the two X dipoles and $\delta\omega_{V,XX_1} = 1.85 \pm 0.05$ GHz and $\delta\omega_{V,XX_2} = 1.66 \pm 0.05$ GHz for the two XX dipoles. In an ideal system, the homogeneous broadening of the linewidth is limited by its lifetime giving a Lorentzian lineshape with a transform-limited linewidth of $\delta\omega = 1/2\pi\tau_b$ ³⁹. We observe Voigt lineshapes indicating a Gaussian component (i.e. inhomogeneous broadening) likely from a fluctuating charge environment⁴⁰ common with above-band excitation. Using measured lifetimes, we calculate transform-limited linewidths of $1/2\pi\tau_{b,X} = 0.0947$ GHz for X and $1/2\pi\tau_{b,XX} = 0.162$ GHz for XX . Observed linewidths are 19.3 and 27.3 (11.4 and 10.2) times this limit for X (XX), depending on the dipole, which gives a measure

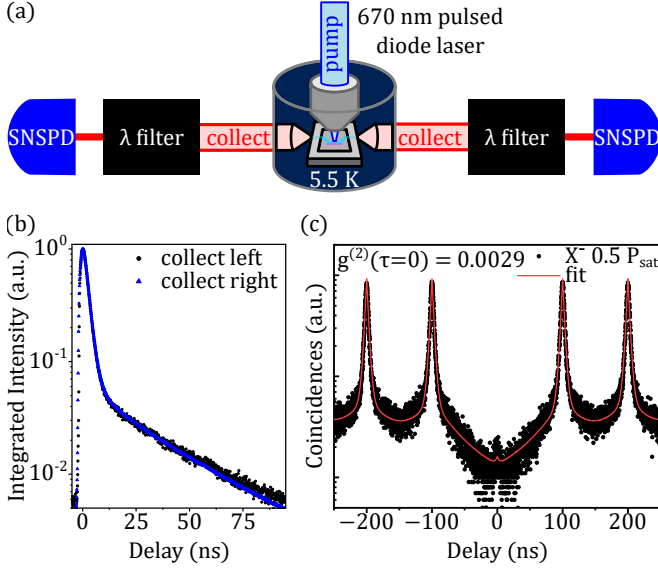


FIG. 6. (a) Schematic of bi-directional collection from the gap-coupled device. (b) Comparison of the PL decay of X^- emission for collection from the left and right facets of the chip. (c) Coincidences, $g^{(2)}(\tau)$, measured using the configuration in (a). For each measurement, the excitation power was $0.5P_{\text{sat}}$.

of the charge noise. The reduced excess broadening observed for XX may be due to the shorter lifetime, τ_b , of this complex, making it less susceptible to charge fluctuations on this time scale.

For the X^- transition, shown in Fig. 5(c), a single peak is observed, consistent with its identification as a singly charged complex and thus not expected to have a FSS^{37} . The lineshape here is also Voigt with a linewidth of $\delta\omega_{V,X^-} = 1.311 \pm 0.007$ GHz, 17.4 times the transform-limited linewidth of $1/2\pi\tau_{b,X^-} = 0.0755$ GHz. We note that above-band excitation is not expected to produce lifetime-limited photons⁴¹, however, the nanowire quantum dot system has been shown to approach this limit³⁹. It is possible that slower Group V switching in this sample may have resulted in more pronounced compositional tails³⁸, effectively increasing the dot thickness and making it more susceptible to fluctuations in the charge environment.

Measurements thus far have demonstrated efficient coupling of single photons on chip, similar to previous experiments using straight waveguides^{24,25}. Next we harness the versatility of the gap-coupled architecture to perform bi-directional measurements. In this case, we excite the nanowire from above using free-space optics and collect from both edge facets, as shown in Fig. 6(a). Fig. 6(b) compares time-resolved PL traces of the X^- emission for collection from the left and right facets. The two traces lie on top of each other with decay times equivalent to that obtained when excitation was via a waveguide, see Fig. 4(f). This measurement simply confirms that the excitation and collection configuration does not

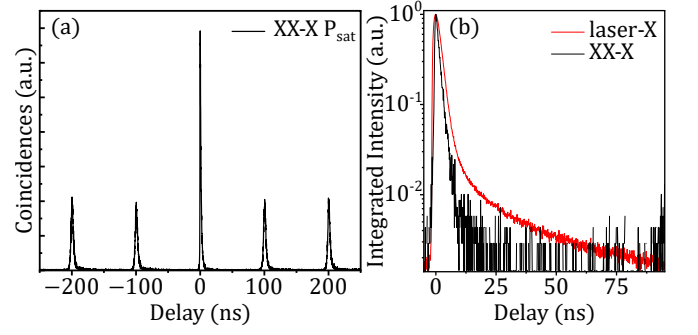


FIG. 7. (a) $XX-X$ cross-correlation measured using bi-directional collection at P_{sat} . (b) PL decay of the X emission triggered using XX detection (black curve) and a synchronization pulse from the laser (red curve).

impact the decay processes of the transition.

Next we perform a second-order correlation measurement where collection from both ends of the nanowire replaces the beamsplitter in the HBT set-up (compare Fig. 4(a)). The $g^{(2)}(\tau)$ measured at an excitation power of $0.5P_{\text{sat}}$ is shown in Fig 6(c). Using Eq. 2, we obtain $g^{(2)}(\tau = 0) \sim 0.003$, increasing to $g^{(2)}(\tau = 0) \sim 0.013$ at P_{sat} . The low level of coincidences in the zero delay peak confirms, again, the single photon nature of the emission from the quantum dot.

Using the same configuration, we can also perform cross-correlation measurements by filtering for different complexes at each output. Fig. 7(a) shows the cross-correlation between XX and X photons where we collect X photons from one side of the chip and XX photons from the other. With XX is sent to the 'start' detector and X to the 'stop' detector we observe strong bunching for $\tau > 0$ consistent with the cascaded $XX-X$ emission process¹⁸.

In Fig. 7(b), we zoom in on the bunched peak at short delay. This time-resolved PL (black curve) represents the decay of the X photon where the start trigger is the detection of an XX photon collected from the opposite facet. For comparison, we also show the decay obtained using a synchronization pulse from the laser (red curve) as in Fig. 4(b). Using the laser trigger, we observe a bi-exponential decay as before with similar extracted parameters: $t_{b,X} = 1.765 \pm 0.003$ ns, $t_{d,X} = 16.2 \pm 0.3$ ns and $A_d/A_b = 0.036$. Triggering on the XX photon, we observe a mono-exponential decay with an extracted time constant of $t_b = 1.442 \pm 0.005$ ns, over 300 ps faster. Since detection of the XX photon indicates that the system is in the bright X state, the slow decay process, associated with a transition from bright to dark via a spin flip at a rate $1/t_{d,X}$, will never be observed. The 300 ps shorter decay extracted from the cross-correlation trace likely represents a more accurate measure of the radiative lifetime since it is not polluted by excitation and spin-flip processes.

Lastly, we determine the coherence time of the X^- photons by measuring the first-order correlation,

$g^{(1)}(\delta\tau_p)$, using an unbalanced Mach-Zehnder interferometer (MZI) with tunable delay, $\delta\tau_p$, and phase modulator, ϕ . The measurement is made in two configurations: a standard setup using an external beamsplitter at the input of the MZI, Fig. 8(a), and a setup that uses the bi-directional nature of the gap-coupled device as a beamsplitter, Fig. 8(c).

The interference fringes, see inset Fig. 8(b), are obtained from the SNSPD counts with the phase modulator set to $\omega/2\pi = 30$ MHz. The interferogram is fit for different delays, $\delta\tau_p$, between the two arms using a sine function, $y = A \sin \omega t + y_0$ and the fringe visibility is calculated as A/y_0 where A is the amplitude of the fringe and y_0 is the y-offset of the count rate. The visibility, plotted in Figs. 8(b) and (d), is highest when the MZI is balanced and decreases with increasing delay between the two arms. In an ideal MZI, one expects visibilities at $\delta\tau_p = 0$ to be 100%. Instead, we observe maximum visibilities of $\sim 73\%$ for the single-sided measurement and $\sim 42\%$ for the double-sided measurement which we attribute to a difference in count rate in the two arms of the interferometer.

For a homogeneously broadened transition, an exponential decay of the fringe visibility with a time constant, T_{hom} is expected. Consistent with the above linewidth for the X^- photon, Fig. 5(c), inhomogeneous broadening gives rise to a Gaussian component resulting in a Voigt profile⁴². To fit the decay in Fig. 8(b), the fringe visibility is appropriately described with a Voigt profile using⁴³

$$g^{(1)}(\delta\tau_p) \sim \exp \left[-\frac{\pi}{2} \left(\frac{\delta\tau_p}{T_{\text{inhom}}} \right)^2 - \frac{|\delta\tau_p|}{T_{\text{hom}}} \right], \quad (3)$$

where T_{inhom} is the Gaussian component of the decay due to inhomogeneous broadening. Extracted T_{hom} and T_{inhom} values are used to calculate the coherence time, τ_c , from⁴⁴

$$\tau_c = T_{\text{inhom}} \exp \left(\left(\frac{T_{\text{inhom}}}{\sqrt{\pi} T_{\text{hom}}} \right)^2 \right) \text{erfc} \left(\frac{T_{\text{inhom}}}{\sqrt{\pi} T_{\text{hom}}} \right) \quad (4)$$

from which we obtain $\tau_c = 180.43$ ps for the standard configuration and $\tau_c = 178.48$ ps with the gap-coupled device acting as the input beamsplitter. To compare with the linewidth measurement, Fig. 5(c), we also calculate the FWHM of the Voigt profile in the frequency domain given by¹⁸

$$\delta\omega_V = 0.535\delta\omega_L + \sqrt{0.217\delta\omega_L^2 + \delta\omega_G^2}, \quad (5)$$

where $\delta\omega_L = (\pi T_{\text{hom}})^{-1}$ is the Lorentzian contribution and $\delta\omega_G = \sqrt{2 \ln 2} / \sqrt{\pi} T_{\text{inhom}}$ is the Gaussian contribution. Extracted coherence times and Voigt linewidths are listed in Table II.

Extracted coherence times and linewidths from the first-order correlation measurements are similar for both

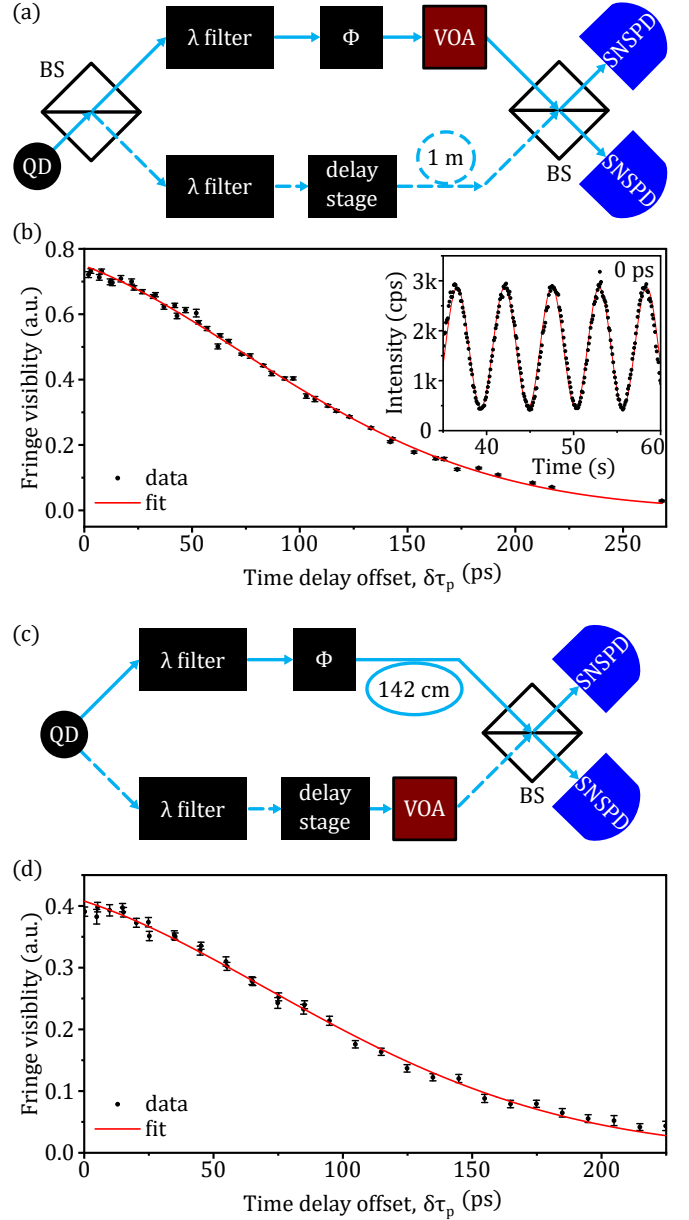


FIG. 8. (a) Schematic of the coherence setup using an external beamsplitter, BS, with collection from one side of device. (b) Fringe visibility as a function of delay, $\delta\tau_p$, between the two arms of the interferometer. Inset shows the interferogram at $\delta\tau_p = 0$. (c) Schematic of the setup with collection from both sides of the device. (d) Fringe visibility measured with the gap-coupled configuration acting as the beamsplitter. The variable optical attenuator (VOA) is used to balance the count rates incident on the two ports of the output beam-splitter.

measurement configurations, as expected. Comparison with values extracted from the linewidth measurements reveal much shorter coherence times in the correlation measurements. Linewidth and $g^{(1)}(\delta\tau_p)$ measurements probe the coherence of the emitted photons separated in time according to the duration of the measurement. In both cases, this time scale is on the order of nanoseconds

and we expect extracted coherence times to be similar. A possible cause of the large discrepancy may be related to imperfections in the interferometer from, for example, the fibre-coupled delay stage not maintaining a constant throughput throughout its tuning range. The discrepancy, however, persisted even after correcting for the interferometer throughput by calibrating using a coherent laser.

Experiment	τ_c (ps)	$\delta\omega_V$ (GHz)
$g^{(1)}(\delta\tau_p)$ single-sided	180.43	3.825
$g^{(1)}(\delta\tau_p)$ double-sided	178.48	3.883
Emission linewidth	1213.81	1.311

TABLE II. Extracted coherence times, τ_c , and linewidths, $\delta\omega_V$, of the X^- photons from different experiments. τ_c is calculated from Eq. 4

V. CONCLUSION

We describe a versatile hybrid platform for efficiently generating single photons on chip based on III-V nanowire quantum dots integrated with Si_3N_4 photonic circuitry. The approach utilizes a nanomanipulator to place individual nanowires on curved Si_3N_4 waveguides providing evanescent coupling of single photons emitted from both ends of the nanowire. We use the bi-directional collection to demonstrate single photon emission without the need of a external beamsplitter. The bi-directional collection is also used to perform cross-correlation measurements for accurate determination of the exciton radiative lifetime as well as first-order correlation measurements to determine the coherence of the emitted photons.

The approach is currently limited by the uncertainty of the dot position within the nanowire (see supplementary material) and by the precision with which the nanowire can be placed with respect to the Si_3N_4 circuitry. Additionally, the measured nanowires were grown with slight tapers whereas the device designs are based on perfectly untapered structures. Further growth studies are expected to lead to better control of the nanowire taper²² and can also address the growth incubation

time⁴⁵, in part responsible for the observed variation in dot position⁴⁶. Improved nanowire placement can be expected with advances in nano-positioning technologies⁴⁷ and their automation will address the scalability of device assembly, limited here by the manual nature of the process.

In this study the sources were operated using above-band excitation, known to cause fluctuations in the charge environment⁴⁸. Such incoherent excitation introduces excess broadening well above the transform limit⁴¹ and limits the indistinguishability of the generated photons⁴. State-of-the-art sources of indistinguishable photons⁵⁻⁹ are operated using resonant excitation⁴⁹ and employ microcavity structures²⁸ to increase radiative decay rates⁵⁰. Further development of the current platform thus requires implementation of resonant excitation, already demonstrated in on-chip approaches^{12,30,51,52}. For enhancing decay rates, microcavities based on ring resonators can be used⁵³, a simple extension of the current curved waveguide architecture.

The platform paves the way for designing scalable architectures that can efficiently harvest $> 90\%$ of the photons emitted from the nanowire quantum dot. Such architectures are useful for the study of coherent light-matter interactions with solid-state two-level systems. They will also be required in future quantum information processing technologies based on the availability of multiple on-chip sources of high coherent single photons.

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DATA AVAILABILITY

Data supporting the findings of this study are available from the corresponding author upon reasonable request.

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Nanowire quantum dot emitters evanescently coupled on-chip via curved and segmented SiN waveguides

Edith Yeung,^{1,2} Kataryna Sorensen,^{1,2} David B. Northeast,² Maziyar
Milanizadeh,² Philip J. Poole,² Robin L. Williams,² and Dan Dalacu^{1,2}

¹*University of Ottawa, Ottawa, Ontario, Canada, K1N 6N5.*

²*National Research Council Canada,
Ottawa, Ontario, Canada, K1A 0R6.*

I. ESTIMATED DOT - WAVEGUIDE COUPLING EFFICIENCY

In Table S1 we present estimates of the losses in the optical system. The system loss, obtained from a throughput measurement using a laser, includes off-chip optical components such as the fiber-coupled tunable filter and associated fiber mating sleeves. The waveguide propagation loss, scattering loss at the window openings, and facet coupling loss are obtained from measurements on a cut-back chip fabricated in the same foundry run. Measured propagation losses were -1.5234 dB/cm for the TE mode and -1.3301 dB/cm for the TM at 975 nm and the length of the waveguide from gap to edge facet was 25 mm. We assume an ideally positioned, completely untapered nanowire with an embedded dot exactly centered along its length such that 50% of the dot emission is directed towards each half of the segmented waveguide. We further assume a coupling efficiency between the dot and the HE_{11} waveguide mode in the nanowire of $\beta_{\text{HE}_{11}} = 95\%$ ¹ and a 20% loss from dot emission into acoustic phonon sidebands² which is filtered-out.

Source of loss	Transmission (%)	Loss (dB)
Detector efficiency	82	-0.862
Optical system loss	10	-10.0
Propagation loss	91.6 (TE)/ 92.6 (TM)	-0.3809 (TE)/ -0.3325 (TM)
Facet loss	86.9 (TE)/ 68.2 (TM)	-0.6115 (TE)/ -1.6609 (TM)
Window loss	79.3 (TE)/ 82.4 (TM)	-1.006 (TE)/ -0.8413 (TM)
Emission direction	50	-3.0
$\beta_{\text{HE}_{11}}$	95	-0.223
Phonon sidebands	80	-0.969
Total	1.97 (TE)/ 1.62 (TM)	-17.0524 (TE)/ -17.8467 (TM)

TABLE S1. Breakdown of losses for an ideal case with the QD within an untapered nanowire exactly centered in the gap between waveguide segments.

Detected counts for X and X^- were 575 kcps and 330 kcps, respectively, for a total count rate of $I_{\text{meas}} = 905$ kcps (i.e. X and X^- are mutually exclusive transitions). Assuming half the emitted photons in the InP nanowire couple to the TE mode of the Si_3N_4 waveguide and half to the TM mode, the total in-fiber count rate is given by $2I_{\text{meas}}/(\eta_{\text{TE}} + \eta_{\text{TM}})$ where $\eta_{\text{TE}} = 0.0197$ and $\eta_{\text{TM}} = 0.0162$. Dividing by the laser repetition rate of 80 MHz, we obtain

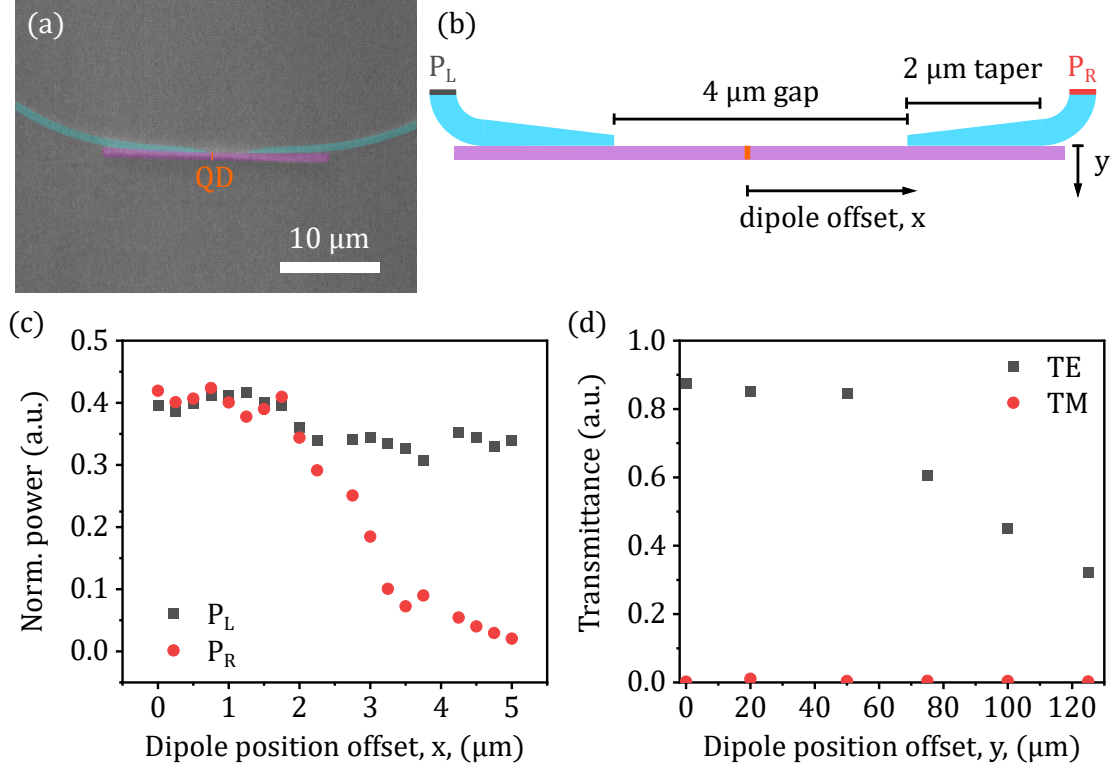


FIG. S1. (a) SEM image of a gap coupler (blue) with a nanowire (purple) where the QD is schematically shown in orange at the nominal position in the middle of the nanowire. (b) Schematic (not to scale) showing the location of ports (P_L and P_R) used for the simulations. (c) Calculated power transfer from the InP nanowire to the P_L and P_R ports as the dipole is shifted from the nominal position, $x = 0$. (d) Calculated transmission of both TE and TM modes through the device as the nanowire is displaced from the Si_3N_4 waveguide from the nominal value of $y = 0$.

an estimated InP nanowire to Si_3N_4 waveguide coupling efficiency of 63% compared to the simulated value of $> 90\%$ from Fig. 2 of the main article.

II. SIMULATIONS FOR POWER TRANSFER DEPENDING ON QD POSITION

The intention of this work's device design is to have the nanowire positioned flush to the Si_3N_4 waveguide tapers and for the quantum dot to be located midway within the gap between the tapers. In this section we investigate the influence of two possible misalignment cases. The first case is axial misalignment of the dipole relative to the gap. The second case looks at instances where the nanowire is increasingly displaced from the Si_3N_4 tapers.

In Fig. S1(a), the dipole position is varied axially in the nanowire, referred to as x in Fig. S1(b), from a centred position ($x = 0\ \mu\text{m}$) to a final dipole position just beyond the tapered waveguide section ($x = 5\ \mu\text{m}$). The dipole is moved only toward the right-side waveguide. Powers in both waveguides, P_L and P_R , are normalized to the emitted dipole power. Fig. S1(a) shows that, while the dipole is located in the $4\ \mu\text{m}$ gap section, roughly equal power is delivered to both silicon nitride waveguides. Once the dipole is beside the tapered region of the waveguide, P_R steadily drops as the adjoining waveguide becomes wider. This is in contrast to P_L , which does not drop considerably when the dipole is located beside the right-hand taper. We suspect that this effect is the largest source of count rate asymmetry in our device and suggests a displacement of several microns to account for the $\sim 100\times$ asymmetry observed.

The second series of simulations shows the influence of increasing displacement of the nanowire away from the waveguide, denoted as y in Fig. S1(b). Considered here is only the HE_{11} mode that is aligned to the TE polarization, but the principle will apply equally to the TM mode. Transmittance of mode power into the waveguide is nearly constant until these displacements are greater than $50\ \text{nm}$. At $125\ \text{nm}$, only $\sim 30\%$ of mode power is transmitted into the waveguide. The insensitivity displayed to displacements below $50\ \text{nm}$ is due in part to the adiabatic coupling provided by the tapered waveguiding sections.

III. QUANTUM DOT LOCATION WITHIN THE NANOWIRE

To determine the location of the quantum dot within the nanowire, we grew a calibration sample where the growth was stopped right after inserting the dot. An SEM image of an array of these nanowire 'cores' is shown in Fig. S2(a). From measurements of 27 nanowires, we obtain a mean height of $11.1\ \mu\text{m}$ with an uncertainty of $\pm 0.7\ \mu\text{m}$. This distribution is promising for targeting dot placement to within a couple of microns of the midpoint of the gap, e.g. to simultaneously maximize counts at both P_L and P_R . It also suggests that the studied device had a core height several standard deviations from the mean if the $\sim 100\times$ count rate asymmetry in P_L and P_R is primarily due to the dot displacement within the gap.

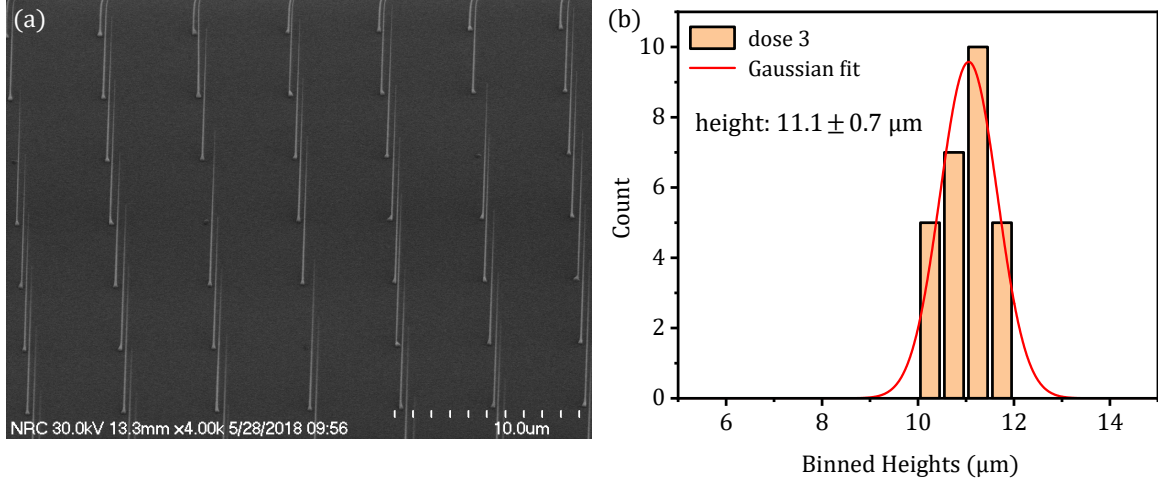


FIG. S2. (a) An SEM image at 45° of an array of nanowire cores defining the position of the quantum dot within the nanowire. (b) Histogram of the heights obtained from measurements of 27 nanowire cores.

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- ¹ I. Friedler, C. Sauvan, J. P. Hugonin, P. Lalanne, J. Claudon, and J. M. Gérard, “Solid-state single photon sources: the nanowire antenna,” *Opt. Exp.* **17**, 2095 (2009).
- ² P. Lodahl, S. Mahmoodian, and S. Stobbe, “Interfacing single photons and single quantum dots with photonic nanostructures,” *Rev. Mod. Phys.* **87**, 347 (2015).

5.2 Conclusion

In this chapter, we have demonstrated a different on-chip architecture for recovering the emission from both ends of a nanowire-quantum dot. The simulations predicted equal collection from both ends for the gap coupler geometry. In practice this was not the case. As stated in the article, we assume this is due to a misplaced position of the dot in the gap. In future experiments, this assumption will be verified by incorporating multiple energy-tuned dots within the nanowire to identify the coupling dependence on dot position [27].

Chapter 6

Plug and Play – Fibre-packaged Hybrid PIC Package

6.1 Introduction

In the previous chapters, we coupled onto our hybrid PIC devices using lensed fibres which required active alignment. Although this setup is useful for quickly testing multiple devices with different architectures under development, we ultimately want to move towards a field-deployable plug-and-play single-photon source. In this chapter, we present a fibre-packaged hybrid PIC chip. The first section will detail the design process of the chip. In the second section, the experimental characterization for a packaged chip will be presented. The results and discussion for multiple cooldown cycles is presented in the paper entitled, “Chip-integrated Photonics-based Fibre-packaged Plug-and-play Single-Photon Sources.”

6.2 Simulations and Design

The goal for this project was to develop a plug-and-play single-photon device using the NW-QDs grown at NRC. Thus far, a key advantage of placing nanowires on chip was to have an emitter in-plane to the circuit with the intention of building more complex circuit designs (such as filters, splitters, gates, and detectors) into the platform in the future.

LioniX International offered a 850 nm multi-project wafer process and an additional packaging service that allowed for fixed alignment of a PIC to a fibre array. As part of the fibre-packaging process, the chip facets are polished and fixed to a mount with fibre arrays. At the time of conceptualization, this was the only wafer process that used this packaging service although it was not the ideal choice for the NW-QDs. This process guided us to using 900 nm emitters.

The chip is designed for NW-QDs emitting at 900 nm. The LioniX process design kit (PDK) specifies that

- the chip dimensions are 1 cm by 1 cm,

- the Si_3N_4 layer is 100 nm thick,
- a 25 nm depth of the waveguide is exposed in the partially clad waveguide regions (windows) (i.e., the NWs can only be placed on top of the waveguide),
- the minimum feature size is 400 nm due to their photolithography process,
- the minimum bend radius is $100\text{ }\mu\text{m}$ due to leaky waveguide modes,
- the fibre array contains a set of twelve polarization maintaining (PM) 780HP fibres with a pitch of $127\text{ }\mu\text{m}$, each aligned to the TE mode of the waveguide, and
- two of the fibres at the top and two at the bottom of the array are coupled for use in the alignment of the fibre array relative to the PIC. This limits us to only eight usable devices.

With these limitations in mind, we need to determine the following:

1. the waveguide width for single-mode operation (i.e., fundamental TE and TM modes) for 900 nm light in a. the fully clad region (see Fig. 6.2 (a)) and b. the window (see Fig. 6.2 (b)),
2. facet width for efficient coupling between the fibre and waveguide facet, and
3. taper lengths to minimize reflections of the tapered couplers at the facets and towards the window region.

A schematic of the different components and geometries of the waveguide to be determined through simulations is shown in Fig. 6.1. The straight facet piece at the edge couplers is added to ensure the correct waveguide geometry even after chip polishing as per LioniX's specifications.

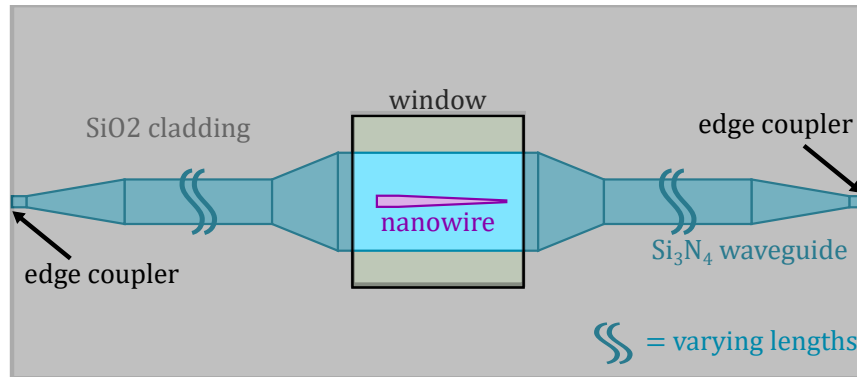


Figure 6.1: Schematic of a waveguide on chip. Not to scale.

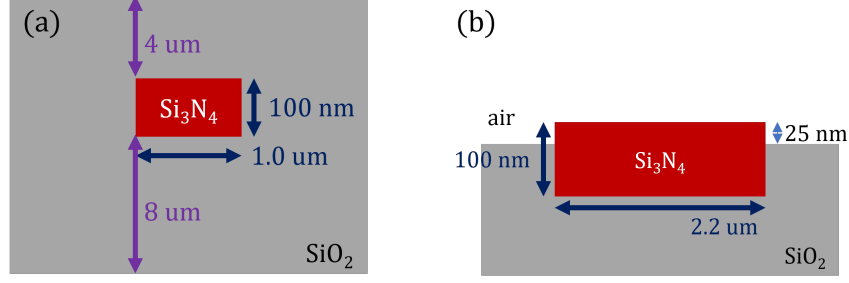


Figure 6.2: Cross sectional schematic of the photonic waveguide chip at (a) a fully clad region and (b) partially clad region where the SiO_2 top cladding is etched down exposing part of the Si_3N_4 .

6.2.1 Waveguide Geometry

First, the waveguide's geometry in the fully clad region is determined as shown in Fig. 6.2 (a).

The refractive indices of Si_3N_4 and SiO_2 at room temperature are modeled with a Cauchy distribution provided for LioniX's TriPlex waveguide materials. This model estimates the refractive index, including birefringence, for the two materials. The refractive index of Si_3N_4 for the TE and TM modes are given by

$$\begin{aligned} n_{\text{Si}_3\text{N}_4(\text{TE})} &= A + B/\lambda^2 + 0.0045, \\ n_{\text{Si}_3\text{N}_4(\text{TM})} &= A + B/\lambda^2 - 0.0045, \end{aligned} \quad (6.1)$$

where $A = 1.97230$ and $B = 0.0161$ are the Cauchy parameters. The wavelength λ is in μm . The refractive index of SiO_2 for the TE and TM modes are given by

$$\begin{aligned} n_{\text{SiO}_2(\text{TE})} &= A + B/\lambda^2 - 0.0004, \\ n_{\text{SiO}_2(\text{TM})} &= A + B/\lambda^2 + 0.0004, \end{aligned} \quad (6.2)$$

where $A = 1.44345$ and $B = 0.00593$.

The electromagnetic field of the modes propagating in the waveguide is mostly confined in the higher index Si_3N_4 waveguide core and evanescently decays in the lower index SiO_2 cladding. Since the field does not stay entirely in one material, the modes experience an effective refractive index between the indices of the two materials in which the field occupies. This is referred to as the effective mode index, n_{eff} . The closer n_{eff} is to the refractive index of the core, $n_{\text{Si}_3\text{N}_4}$, the more strongly confined the guided mode [103].

The effective mode index, n_{eff} , is obtained by deconstructing a rectangular waveguide into two one-dimensional problems (similar to a slab waveguide). The thin one-dimensional slab waveguide is analyzed to solve for the propagation coefficient, β , for the wavelength and mode of interest. The effective mode index is determined through

$$n_{\text{eff}} = \frac{\beta}{k_0}, \quad (6.3)$$

where β is the propagation coefficient and k_0 is the vacuum wave-vector of the light being guided [104].

The effective modes indices of propagating modes were calculated using the Wave Optics module in COMSOL given a fixed height of 100 nm for the Si_3N_4 core and taking a parameter sweep of the waveguide width (0.25 μm to 3 μm). The simulation also sweeps across different propagating wavelengths: 850 nm (the intended wavelength specified by LioniX), as well as 900 nm and 950 nm, which are relevant for NW-QD emission. For each set of parameters, the search method solves for mode indices at the near-bulk value and smaller. The found modes should include the fundamental TE and TM modes.

In the simulation for the fully clad waveguide, the surrounding SiO_2 cladding was set to have a height and width of 10 μm , which was sufficient for modelling the mode within the Si_3N_4 core with a negligible effect from the edges of the SiO_2 box. Fig. 6.3 (left column) shows the modes present for a fully clad waveguide.

LioniX offers a window etching process that allows for the Si_3N_4 waveguide layer to be exposed in a region of choice. Their process etches the top cladding such that 25 nm (from the top) of the waveguide is exposed. These window regions allow us to place the NW-QDs for evanescent coupling to the nitride waveguides. On chip, these regions were 100 μm by 100 μm wide. We also investigated the waveguide modes in these regions to determine how disparate they are from the fully clad regions. For the simulation, we set the top portion to vacuum and the bottom to SiO_2 such as in Fig. 6.2 (b). The results for the same parameters sweeps for a system with this geometry is shown in Fig. 6.3 (right column). The NW-QD will be placed on the waveguide in the exposed region.

From the two sets of simulations (fully and partially clad), we want to select the widths that will contain only the fundamental TE and TM modes (TE_0 and TM_0) for our selected range of wavelengths while also minimizing the difference in widths between the clad and partially clad regions. We also want an n_{eff} closer to the refractive index of the core for better confinement in the core. It can be seen in Fig. 6.3 that as the wavelength increases, n_{eff} decreases. Most importantly, we want widths in each region that only support the fundamental waveguide modes. The width of the waveguide is chosen in the fully clad region to be 1.0 μm , so that it would be at the single-mode cutoff for 850 nm.

Although the fundamental TE and TM modes are always present in a symmetric waveguide, this is not the case for an asymmetric waveguide in the window region [103]. Notice in Fig. 6.3 (right column), there is no TM_0 mode in the waveguide width range simulated. The electric field in the TE modes is mostly parallel to the substrate surface, whereas in the TM modes, the electric field has a strong vertical component which interacts more strongly with the asymmetric interface. When the waveguide is asymmetric, the field distribution shifts toward the higher-index side, in this case, downward toward the SiO_2 substrate (see attached paper Ch. 6.4 Fig.1(d-e)). The effective mode index, n_{eff} , decreases compared to a fully symmetric Si_3N_4 waveguide, because part of the field now resides in air, which has a lower index.

There is a drastic difference in the where the modes start to exist as a function of width between the two regions. No TE mode is expected to exist in the window region for the standard single-mode widths in the fully clad region. Therefore changes to the waveguide

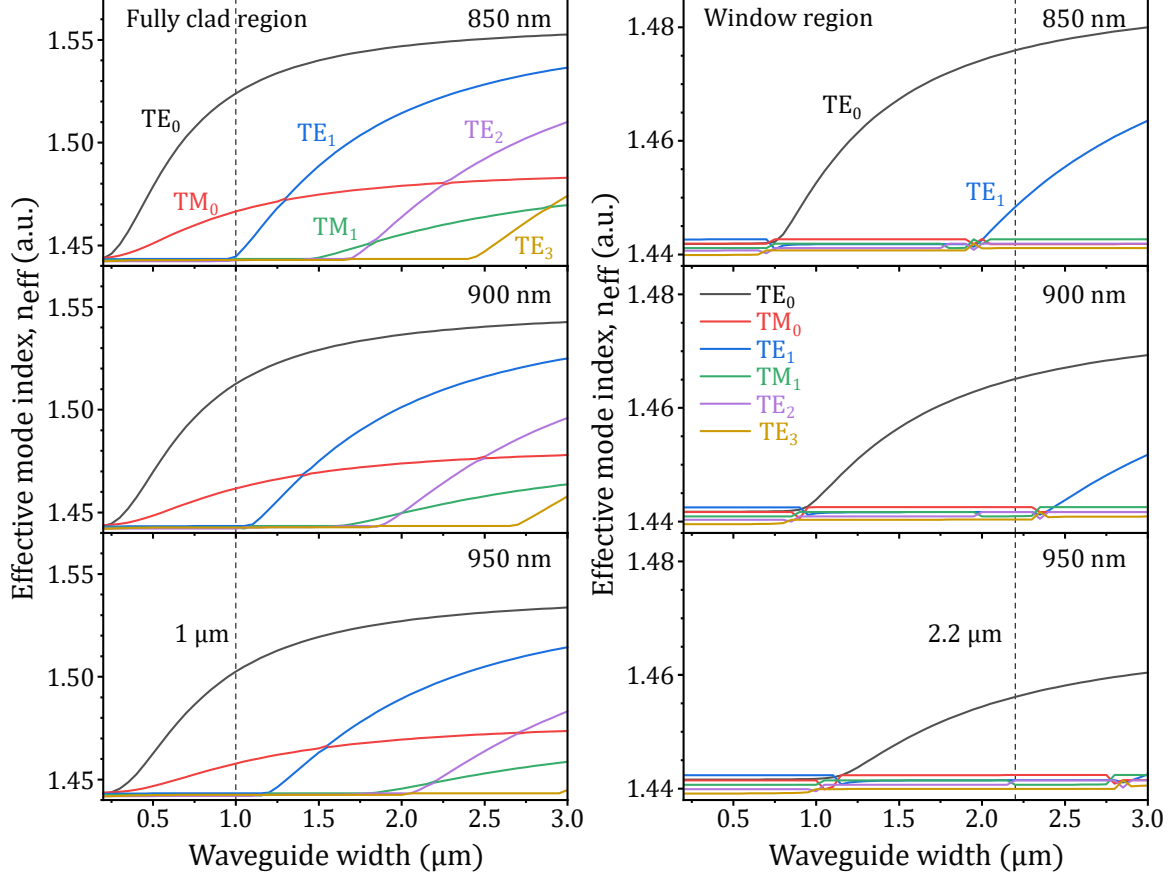


Figure 6.3: Effective mode index as a function of the waveguide width for a 100 nm tall waveguide in (left) a fully clad waveguide and (right) a partially clad waveguide etched down to expose 25 nm of the Si_3N_4 waveguide to air (neglecting the NW-QD). The same legend is used for all six graphs. Note the absence of TM modes in the window region. The chosen waveguide width for the final design is shown with a vertical dashed line.

width in the window region are necessary. A width of $2.2 \mu\text{m}$ was chosen as a compromise between keeping the width down and increasing the index contrast. This ensured that at the emission wavelength of the QD (890~950 nm), the waveguide would only support the TE_0 mode. As demonstrated in Chapter 4, this would prime the system ready for HOM measurements. However, a key disadvantage of a waveguide that only supports one of the two fundamental modes is that half the photons are lost. This is because the dot emits into both the two HE_{11} polarizations that the NW-QD supports, but half of these emitted photon events are lost.

Another reason against designing the waveguide to be too wide relative to the diameter of the NW, is that for optimal power transfer of the NW-QD mode into the underlying waveguide mode, the NW-QD needs to be placed perfectly centred. Using COMSOL's Wave Optics module, we solve for the transmittance, reflectance, and absorptance of the relevant (polarization-aligned) HE_{11} mode of the nanowire to the TE_0 mode of the waveguide for a nanowire ($r_{\text{NW}} = 125 \text{ nm}$, $l_{\text{taper}} = 10 \mu\text{m}$, emitting at 920 nm) placed on a $2.2 \mu\text{m}$

wide waveguide (a) centred, (b) offset by $2r_{\text{NW}}$ from the edge of the waveguide, and (c) placed at a 12° offset across the waveguide. The results are shown in Fig. 6.16 (FIG. S1) in Supplementary Materials of this chapter. For the nanowire placed at an angle, the taper is set to $12\text{ }\mu\text{m}$, allowing for part of it to hang off the side of the waveguide while maintaining a $10\text{ }\mu\text{m}$ length on the waveguide.

In Fig. 6.16 (FIG. S1) (a) in Supplementary Materials, the nanowire placed at the centre showed the highest transmittance. The majority of the power is lost as radiation into the environment consisting of air (above) and the SiO_2 (below). With the nanowire offset towards the edge of the waveguide (see Fig. 6.16 (FIG. S1) (b)), less of the field injected into the system from the quantum dot makes it to the port on the waveguide. Since waveguide only supports the TE_0 mode, the worsened spatial matching between the mode of the nanowire and waveguide caused by the offset lowers the transmittance for that system. A nanowire positioned at an angle (see Fig. 6.16 (FIG. S1) (c)) also suffers from reduced spatial matching of the modes compared to the ideal case where the nanowire is centred. However, the angle may have contributed to better spatial matching than a nanowire completely offset towards the edge of the waveguide since the highest electric field intensity along the nanowire remains near the centre of the waveguide.

It will be difficult to place a nanowire perfectly centred along the waveguide, so limiting the waveguide width to $2.2\text{ }\mu\text{m}$ is simply more practical. This also limits the length of taper needed for the transition between the two widths. Further discussion on the taper lengths will be presented in Ch. 6.2.3.

6.2.2 Edge coupler facet width

For optimal coupling between the fibre and the waveguide, the mode in the fibre and mode in the waveguide need to be matched at the interface. In practice, the waveguide facet is designed to maximize the modal overlap integral with the optical fibre mode. The Finite Difference Eigenmode (FDE) solver in Lumerical is used to calculate the spatial profile and frequency dependence of modes by solving Maxwell's equations on a cross-sectional mesh of the fibre or waveguide to find the 2D spatial field profiles, effective index, and estimated loss.

To find the optimal facet width, we first find the mode profile for a 780 HP fibre with a core diameter of $4.4\text{ }\mu\text{m}$. The material indices for Si_3N_4 [105] and SiO_2 [106] were taken from literature [107]. In Fig. 6.4 (a), since the expected mode field diameter of the fibre mode is $5.0 \pm 0.5\text{ }\mu\text{m}$ at 850 nm [108], the simulation region was restricted to a $20\text{ }\mu\text{m}$ by $20\text{ }\mu\text{m}$ rectangular subregion centred around the core. It can be seen in the resulting mode profile that this size is sufficient for calculating the entire mode around the core.

The electric field profile for a 250 nm by 100 nm Si_3N_4 waveguide is shown in Fig. 6.4 (b). The simulation accounts for a 500 nm thick layer of gold sitting at the top of the cladding. The top gold layer is used by LioniX for labels and alignment marks (see photograph in Fig. 6.7 (a)). From the previous simulation of the eigenmodes present at each waveguide width (see Fig. 6.3 (top left)), we want to limit the facet width to less than the single-mode cutoff at $1\text{ }\mu\text{m}$ for 850 nm .

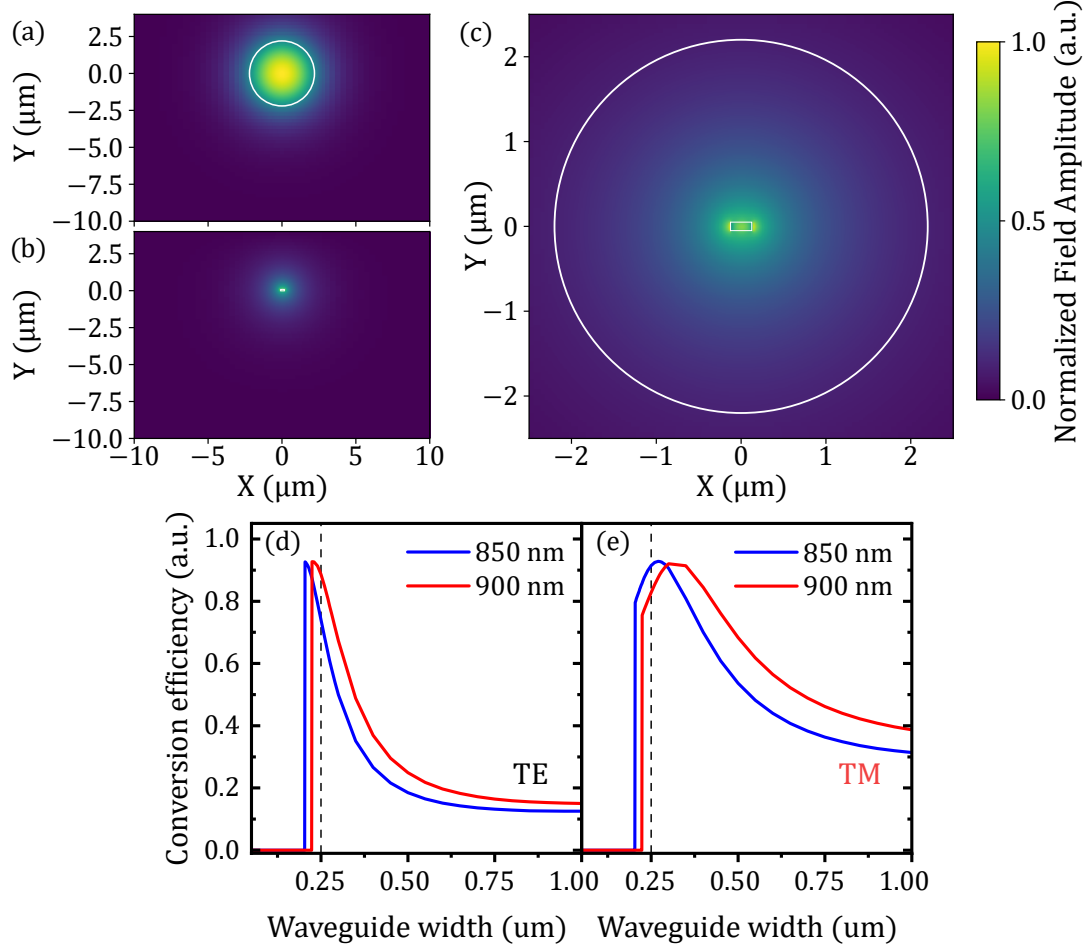


Figure 6.4: Mode profiles are shown for (a) a 780HP fibre with a 4.4μ diameter core and (b) a 250 nm by 100 nm Si₃N₄ waveguide in SiO₂ cladding. The outline of the fibre core and waveguide are shown in white. A magnified version is shown in (c) of the waveguide mode profile with both the fibre core and Si₃N₄ waveguide outlines overlaid for clarity. Mode conversion efficiency is plotted for a 780HP fibre and Si₃N₄ waveguide by varying the waveguide width for (d) TE₀ and (e) TM₀ optical modes.

After solving for the modes at the fibre and waveguide facet, Lumerical calculates the field overlap and power coupling between an input mode from the fibre and all of the modes of the facet waveguide. Fig. 6.4 (c) shows a zoom in of the waveguide mode, with both the Si₃N₄ waveguide and facet borders overlaid. From previous simulations, we know n_{eff} is much closer to the n_{clad} than n_{core} when the waveguide is narrower. This fact is reflected in Fig. 6.4 (b,c), in which the electric field amplitude is seen to have expanded in size and now mostly resides in the cladding. The resulting overlaps for the TE₀ and TM₀ modes are shown in Fig. 6.4 (d,e).

We see a sharp rise then quick decay in the overlap in the TE₀ mode at a waveguide width of 205 nm for $\lambda = 850$ nm propagating light and the same at 225 nm for $\lambda = 900$ nm light. The maximum coupling efficiency for the TM₀ mode occurs at a waveguide width

of 270 nm for $\lambda = 850$ nm light and 300 nm for $\lambda = 900$ nm light. The coupling efficiency decays as the waveguide becomes wider than the optimal width because the mode gets more confined in the Si_3N_4 core, i.e., the overlap region is smaller. As the width is decreased from the optimal value, the mode around the waveguide into the cladding becomes larger, making the overlap region restricted by the size of the fibre mode. As the overlap between the two modes become more mismatched, the conversion efficiency is expected to drop off. The sharp cutoff seen in both graphs is likely an artifact of the simulation. The simulation was unable to find any modes that supported a waveguide with a width less than 220 nm.

A waveguide width of 250 nm is chosen for the facet. At this width, the coupling efficiency is 74.0% ($\lambda = 850$ nm) and 88.5% ($\lambda = 900$ nm) for the TE_0 mode, and 91.5% ($\lambda = 850$ nm) and 82.8% ($\lambda = 900$ nm) for the TM_0 mode. From the simulations of n_{eff} in the window region, we know that the TM_0 mode does not exist, so coupling efficiency of that mode is less of a priority. A width of 250 nm was chosen rather than at the optimal position of 225 nm ($\lambda = 900$ nm) to account for fabrication related feature size uncertainties. The NW-QDs chosen for pick-and-place have emission wavelengths around 900 nm.

6.2.3 Waveguide taper lengths

Since the waveguide at the facet, the fully clad region, and the partially clad window all have different widths, an appropriate tapering length needs to be found between each transition. Note that since the TM mode is not guided in the window region, the tapering had to be in the fully clad region. In an adiabatic taper, the rate of change of the waveguide geometry must be slow enough that the mode can adapt without significant reflection or scattering into higher order guided modes and radiation modes.

We use Ansys Lumerical's bidirectional eigenmode expansion (EME) solver to simulate the light propagating through a taper. For the simulation, a tapered waveguide is divided into multiple cells, and the eigenmodes are solved at the interface between adjacent cells. Scattering matrices for each section are then formulated by matching the tangential E and H fields at the cell boundaries. The S-matrix of each cell can be propagated bidirectionally to calculate the S-matrix for the entire device [109]. The S-matrix index corresponds a mode at a port, where port 1 is on the narrow end of the taper and port 2 is on the wider end. The S-matrix is constructed based on a combination of the input and output ports and their corresponding modes. Specifically, we are interested in the transmission at the wide port when launching with either the TE_0 or TM_0 modes from the narrow port. A different S-matrix is calculated for different propagation distances (i.e. taper length).

First, we look at the S-matrix for a waveguide at the facet ($w_{\text{facet}} = 0.25 \mu\text{m}$) tapering up to a width of $w_{\text{wg1}} = 1.0 \mu\text{m}$ in the fully cladded region. As we can see from Fig. 6.3 (left), only two modes are present for the two waveguide widths, so there only needs to be four S-matrix indices as shown in Table 6.1.

S-matrix index	Port	Mode	Sx1	Sx2	Sx3	Sx4
1	port 1 (@ facet)	mode 1 (TE ₀)	4.12e-22	1.40e-29	0.996	6.63e-25
2	port 1 (@ facet)	mode 2 (TM ₀)	6.83e-31	1.85e-20	8.77e-25	1.00
3	port 2 (@ wg1)	mode 1 (TE ₀)	0.996	8.04e-25	9.74e-23	2.83e-31
4	port 2 (@ wg1)	mode 2 (TM ₀)	7.13e-25	0.997	1.14e-30	1.86e-20

Table 6.1: Table shows S-matrix indices with related ports and modes. Sample S-matrix for a 200 μm long taper with 900 nm light propagating through.

Within the S-matrix, the Sxy notation corresponds to an output at y for an input at x. A sample computed S-matrix is shown as a heatmap to the right of Table 6.1. For example, S11 corresponds to the reflection coefficient at output port 1 mode 1 (TE₀) from input port 1 mode 1 (TE₀) whereas S31 corresponds to the transmission coefficient at output port 2 mode 1 (TE₀) from input port 1 mode 1 (TE₀). Since the matrix contains complex scattering parameters, the absolute value squared for the S parameters corresponding to (a) S31: launching into port 1 with TE₀ and collecting TE₀ at port 2, and (b) S42: launching into port 1 with TM₀ mode and collecting the TM₀ mode at port 2 are shown in Fig. 6.5. Here, we are calling it a mode conversion efficiency because we are looking at how much of the power at the input mode at port 1 (narrow end) is transmitted to the second port (wide end).

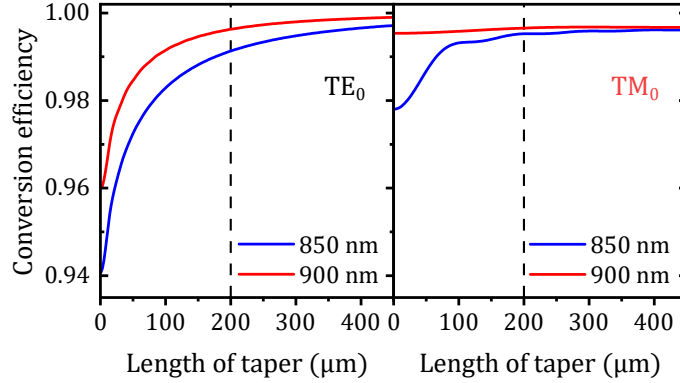


Figure 6.5: Mode conversion efficiency of light propagating through a tapered waveguide with an input port of dimensions 0.25 μm by 0.1 μm to an output port of dimension 1 μm by 0.1 μm . The graph shows the dependence of the mode conversion efficiency as a function of the taper length for light propagating in the (a) TE₀ mode and (b) the TM₀ mode. The final taper length is shown with a vertical dashed line at 200 μm .

For each mode, the simulation was run for 850 nm (blue) and 900 nm (red). It can be seen that the mode conversion efficiency increased for the longer wavelength. We chose the taper length to be 200 μm since it has a conversion efficiency of 99.1% for 850 nm and 99.6% for 900 nm for the TE₀ mode and 99.5% for 850 nm and 99.6% for 900 nm for the

TM₀ mode. In other words, there is no real benefit from making the taper too long since the gains are marginal.

Next, the taper length needed for a waveguide going from $w_{wg1} = 1\ \mu\text{m}$ to a $w_{wg2} = 2.2\ \mu\text{m}$ is determined. We know from Fig. 6.3 (left), that at $2.2\ \mu\text{m}$, there are at least five modes present for the three wavelengths. From top to bottom, they are TE₀, TE₁, TM₀, TE₂, TM₁. This means there needs to be more S-matrix indices to account for the extra modes. Table 6.2 lists the ports and modes to be studied for the S-matrix.

S-index	Port	Mode		Sx1	Sx2	Sx3	Sx4	Sx5	Sx6	Sx7	Sx8
1	port 1 (@ wg1)	mode 1 (TE ₀)	S1y	5.04e-10	2.30e-29	1.48e-30	3.86e-13	1.00	1.69e-26	5.19e-27	8.12e-06
2	port 1 (@ wg1)	mode 2 (TM ₀)	S2y	1.95e-32	3.67e-10	5.98e-29	6.17e-30	5.10e-27	2.67e-21	1.00	1.04e-25
3	port 1 (@ wg1)	mode 3 (TE ₁)	S3y	1.03e-29	1.11e-29	4.67e-11	2.16e-32	1.42e-24	1.00	2.68e-21	8.29e-27
4	port 1 (@ wg1)	mode 4 (n/a)	S4y	3.46e-13	3.52e-29	1.75e-29	1.27e-13	3.32e-06	1.05e-19	7.57e-26	2.70e-11
5	port 2 (@ wg2)	mode 1 (TE ₀)	S5y	1.00	5.97e-27	2.03e-24	3.30e-06	5.08e-10	6.16e-30	2.10e-32	6.26e-11
6	port 2 (@ wg2)	mode 2 (TE ₁)	S6y	1.67e-26	2.69e-21	0.999	8.27e-20	2.99e-30	4.67e-11	8.65e-29	5.26e-33
7	port 2 (@ wg2)	mode 3 (TM ₀)	S7y	6.16e-27	0.999	2.65e-21	9.99e-26	3.10e-29	8.97e-30	3.67e-10	3.19e-32
8	port 2 (@ wg2)	mode 4 (TE ₂)	S8y	8.19e-06	1.04e-25	8.60e-27	2.71e-11	6.15e-11	7.56e-30	2.67e-31	9.59e-10

Table 6.2: S-matrix indices and related port and modes for a tapered waveguide. Note for S3, there is only a hint of the TE₁ mode for 850 nm light in the $1\ \mu\text{m}$ by $100\ \text{nm}$ waveguide. Regardless, this mode is not of interest. The extra modes were included in the simulation because for 850 nm light propagating through a $2.2\ \mu\text{m}$ wide fully-clad waveguide, the TE₂ and TM₀ modes are close to the crossover point. The simulation merely lists solutions (n_{eff}) from largest to smallest. A sample S-matrix is shown on the right.

An example of what the S-matrix looks like for a system where the taper is $50\ \mu\text{m}$ long with $\lambda = 900\ \text{nm}$ is shown to the right of Table 6.2. The matrix shows little mode conversion from one mode input to another. For example, there is negligible rotation of an input TE₀ mode into the TM₀ mode at output S71. There is also little scattering into higher order TE modes at outputs S81 and S61. However, there is strong conversion when the input and output modes at the two ports are the same (S51, S72, S63).

From Table 6.2, it can be seen that S51 and S72 give the transmission coefficients for the fundamental TE and TM modes at the wide output (port 2) from the narrow input end (port 1). Fig. 6.6 shows the conversion efficiency for (a) the TE₀ mode and (b) the TM₀ mode for both 850 nm (blue) and 900 nm (red) light propagating along the taper as a function of the taper length.

A taper of $50\ \mu\text{m}$ is chosen for the waveguide widening from $w_{wg1} = 1.0\ \mu\text{m}$ to $w_{wg2} = 2.2\ \mu\text{m}$ before the window. Again, this length was chosen because there is little benefit in making the taper too long yet still long enough to get high transmission at the output of the taper. This length is sufficient for obtaining a high conversion efficiency of both the TE₀ and TM₀ modes. A schematic of the final waveguide structure is shown in Fig. 6.1.

Using the chosen geometries, the chip design was laid out using the Python package Phidl. The chip contained two sets of waveguides (with and without windows), each with

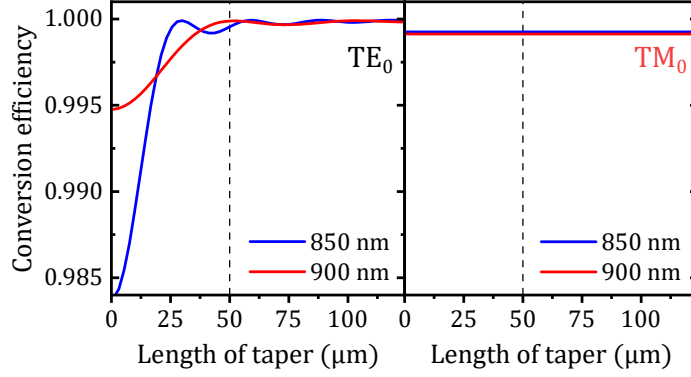


Figure 6.6: Mode conversion efficiency from a waveguide of width $w_{wg1} = 1 \mu m$ expanded to a width of $w_{wg2} = 2.2 \mu m$ as a function of taper length for a propagating (right) TE_0 mode and (left) TM_0 mode. The chosen taper length for the final design is shown with a dashed line at $50 \mu m$.

varying lengths for measuring the propagation loss through the waveguides. An optical image of the final chip is shown in Fig. 6.7 (a). The four loops for fibre alignment are seen on either side of the chip surrounding the middle region. All eight waveguides within this region have windows for nanowire placement – two of which are in large coils for propagation loss measurements. On the bottom half of the chip are four waveguides without windows. The embellishments on the chip were added using the Nazca package. The design of the gold border was provided by LioniX for use as alignment marks in their fabrication process.

6.3 Chip characterization

Four sets of chips were received from LioniX International half a year after the submission of the layout file (Graphical Data System *.gds) for fabrication and packaging. Two of the chips are packaged as seen in Fig. 6.7 (b) and the other two are bare as seen in Fig. 6.7 (a). The latter two chips are used to measure the propagation loss through the system. Only one of the two packaged chips is used for nanowire transfer.

Preliminary characterization include transmission measurements through all accessible waveguides across the chips. The transmission through the waveguides is again measured post transfer to obtain an estimate of the scattering loss due to the nanowire. Each finished device (i.e., fibre-packaged waveguide with nanowire) is characterized for its photoluminescence spectrum and single-photon purity.

6.3.1 Transmission measurements

Transmission measurements are performed to find the propagation loss and facet loss of a set of waveguides from the same batch during fabrication. These measurements are also important for determining which waveguides are worth transferring nanowires to.

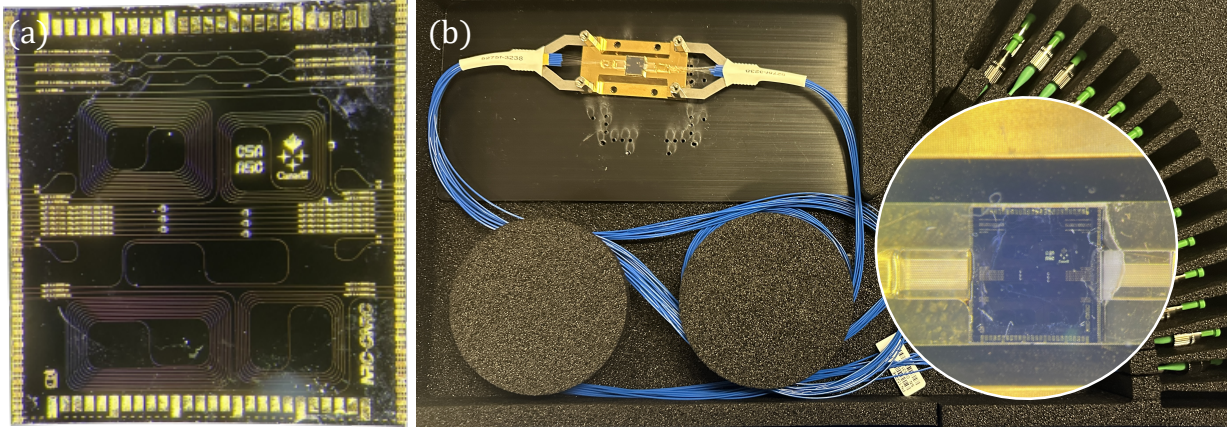


Figure 6.7: (a) Optical image of a 1 cm by 1 cm bare chip with the final waveguide design. Windows can be seen near the centre of the chip as six small boxes staggered blue boxes. (b) Fibre-packaged chip. The 1 cm by 1 cm chip on a 5 cm by 3 cm gold-plated mount. Insert shows a zoom in of the fibre-array on either sides of the chip.

A tunable laser is used to scan across the waveguide from 910 nm to 980 nm. A half-wave plate is used to control the polarization of the laser light launched into the waveguide. The transmission through the waveguide is collected by a fibre on the opposite end of the chip. The light passes through a polarized beam splitter and directed to two detectors.

Initial measurements on a bare chip were performed using lensed fibres and position stages. This method allows access to both the waveguides with and without a window opening. To calibrate the measurement, the transmitted power through the system is found by coupling two lensed fibres to each other as detailed in Ch. 3.1. This measurement (in dBm) is subtracted from the measured transmitted output with the chip to obtain the loss contributions only due to the device under test.

There are four waveguides of lengths $100445 \mu\text{m}$, $50672 \mu\text{m}$, and $20442 \mu\text{m}$, and $9682 \mu\text{m}$ (see coils in the bottom half of the chip in Fig. 6.7 (a)); and six waveguides with window openings. For these latter waveguides there are four repeats of $9682 \mu\text{m}$ and one repeat each of $50072 \mu\text{m}$, and $100234 \mu\text{m}$ (see coils in upper middle portion of chip). A calibrated loss measurement for a set of waveguides without windows is shown in Fig. 6.8 (a) for light launched into the TE mode of the waveguide. One of the waveguides from this set is damaged, but the three remaining waveguides is sufficient for finding the propagation loss. To extract the loss at a specific wavelength, the transmitted power spectrum is fitted with a fifth-order polynomial.

The transmittance as a function of waveguide length is shown in Fig. 6.8 (b) for the set without windows and in Fig. 6.8 (c) for the set with windows. A linear fit of the loss as a function of the waveguide wavelength gives the propagation loss (slope) and facet losses (y-intercept). The loss per facet is assumed to be equal and, so, is half the value of the y-intercept. In each set, the transmittance decreases with increasing wavelengths, consistent with the decreased spatial confinement of the modes in the waveguide, as predicted in the simulations shown in Fig. 6.3. Not shown here are measurements launching into the TE

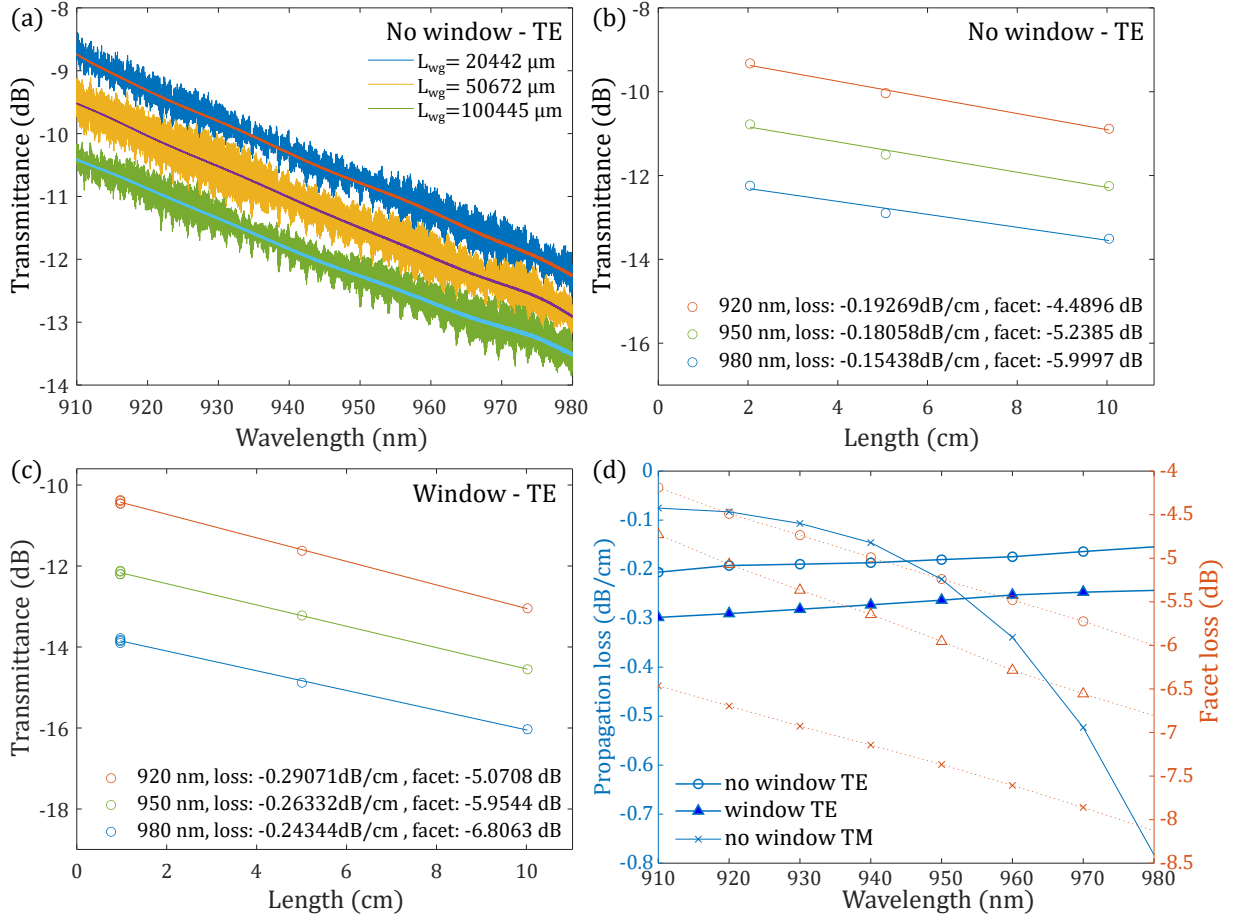


Figure 6.8: (a) Loss measurements for a set of waveguides using lensed fibres. Due to the noise, each plot has been fitted with a fifth-order polynomial. Using the values obtained from the fit of the loss measurements, the wavelength specific loss value is extracted to obtain (b) a transmittance as a function of length. (c) Wavelength specific transmittance as a function of length for a set of waveguides with windows. A linear fit for each wavelength yields the slope which gives the propagation loss and the y-intercept which gives the facet loss. All three of the above was measured by launching light into the waveguide in the TE mode and collecting only the TE mode through a polarized beam splitter. (d) Compilation of the extracted values for the propagation loss (blue) and facet loss (orange) as a function of the wavelength.

and TM modes for both sets of waveguides (i.e., with and without window). It should be noted that for the set of waveguides with a window, when launching into the TM mode, the transmission was negligible (i.e., ~ -40 dB) consistent with simulations which predicted no TM modes propagating through the window regions.

The extracted propagation and facet losses for each set of waveguides as a function of the input wavelength is shown in Fig. 6.8 (d). For the TM mode, the transmittance and the propagation loss drop quickly as a function of wavelength, consistent with the simulations shown in Fig. 6.3.

When launching into the TE mode, when comparing window versus no window, the propagation and facet losses are ~ -0.1 dB/cm and ~ -1.7 dB lower, respectively, at each input wavelength. This difference may be attributed to the scattering losses introduced by the window. The “facet” loss here is a combination of the losses at the facets of the chip and edges of the window (i.e., each window edge contributes ~ -1.7 dB loss). Only the scattering loss due to the window edges is of relevance for this work because the loss at each facet is dependent on the coupling method (i.e., with lensed fibres, cleaved fibres, or fibre array). The additional propagation loss here is due to the difference in the medium of the cladding (i.e., air and SiO_2).

The propagation loss is expected to be the same across the chips regardless of coupling method. The same set of measurements is performed on the fibre-packaged chip to extract the facet loss. Fig. 6.9 shows the loss measurements of a fibre-packaged chip. It becomes immediately obvious that the losses due to the facets are much lower than when coupling with the lensed fibres. However, wavelength-dependent oscillations can also be seen which are not present when using lensed fibres. The oscillations are likely due to a Fabry-Pérot effect caused by the space between the fibre array and facet. As such, it is difficult to fit the curves and extract the propagation and facet losses for all wavelengths. Nevertheless, to smooth out the oscillations, a fifth-order polynomial was used to fit the data (shown in dark hue)

Using the fits from Fig. 6.9 (a), the transmittance from each waveguide was extracted for each wavelength. As before, a linear fit was used to extract the propagation and facet losses. The results are shown in Fig. 6.9 (c). The propagation loss between the two sets (i.e., measured with lensed fibres versus the fibre array) is similar at lower wavelengths (910-950 nm). As seen Fig. 6.9 (a,c), the transmission behaves differently at wavelengths beyond 950 nm. It is unknown why this behaviour occurs in this particular fibre-packaged chip. Not shown here is another fibre-packaged chip that has large oscillations in the transmission across the entire range of wavelengths. The second chip was not chosen for nanowire transfer because the coupling of the fibre-array on one side is suspected to be slightly offset at an angle, leading to the strong oscillations in the transmission data, including through the two alignment loops in the left fibre array.

Since the nanowires are grown to emit around 900 nm, only the lower wavelength data is of significance. Due to the edge couplers being optimized for cleaved fibre coupling, the facet loss shown in Fig. 6.9 (c) (red), show over -2 dB improvement compared to the case with lensed fibres.

The entire fibre-packaged chip was placed on a nanomanipulator stage for nanowire

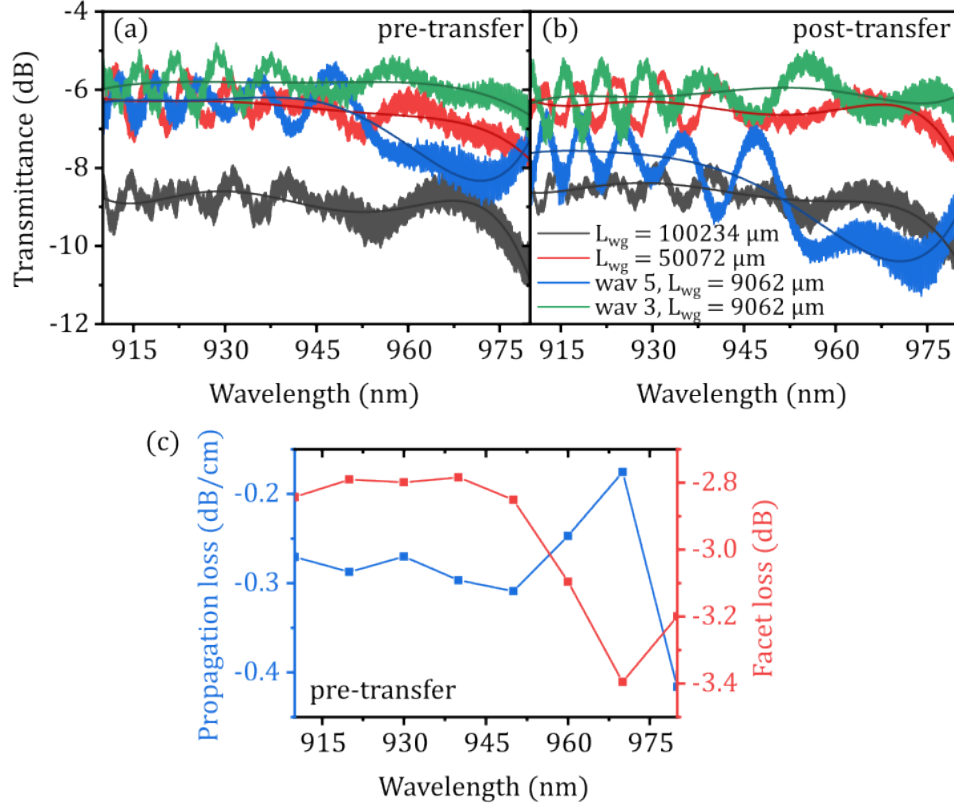


Figure 6.9: Loss measurements for a fibre-packaged chip for waveguides of different lengths both (a) before nanowire transfer and (b) post nanowire transfer. Corresponding fits are shown in a darker hue. (c) The extracted propagation loss (blue) and facet loss (red) for the chip pre-transfer at room temperature.

transfer inside an SEM. To check whether the nanowire transfer process caused any changes to occur to the coupling from the chip to the fibre arrays or to the waveguides, the loss measurements were performed again. The results are shown in Fig. 6.9 (b). Note here that nanowires were only transferred to the four straight waveguides ($L_{wg} = 9062 \mu m$), denoted in the figure as “wav 3” and “wav 5”. The transmission data from the two longer waveguides (red and black) were similar to before transfer, which indicates that the fibre-array and waveguides were not damaged during the transfer process. The transmission through wav 3 (green) was similar to before transfer. However, wav 5 (blue) showed -2 dB lower transmission after the transfer. This may be because there was more stress on the nanowire during transfer to wav 3 which resulted part of the nanowire breaking off.

To measure the sample at cryogenic temperatures, the sample fibres were cut and spliced to a set of fibres inside a Photon Spot Cryostat 4 cryostat. These cryostat fibres were passed through feedthroughs to allow for room temperature access. All fibres were bundled and secured to the 4 K stage of the cryostat. With this configuration, four of the waveguides were accessible without active alignment by simply switching between different sets of fibre connections outside the cryostat.

During the cool down to 4 K, the power throughput of the system was monitored.

Oscillations in the transmitted power (at 930 nm) were observed as the chip cooled, but the power transmitted through the chip stabilized again when it reached base temperature.

6.3.2 Photoluminescence

After the cryostat was cooled, photoluminescence measurements were taken using all accessible fibre ports. On one waveguide, a transferred nanowire had its tapered end break off, leaving just the NW-QD base and the QD intact (see Fig. 6.10 (a)). Using above-band excitation with a 80 MHz 780 nm pulsed laser, nearly equal emission rates from the QD were collected from either ends of the nanowire. Fig. 6.10 (b, c) show a pump and collection scheme where the excitation laser is sent through the waveguide and both the emission from the QD and laser are collected at the other end. Next, we use a 90-10 beam splitter to pump and collect the emission of the QD from one side of the waveguide (see Fig. 6.10 (d) and (e)).

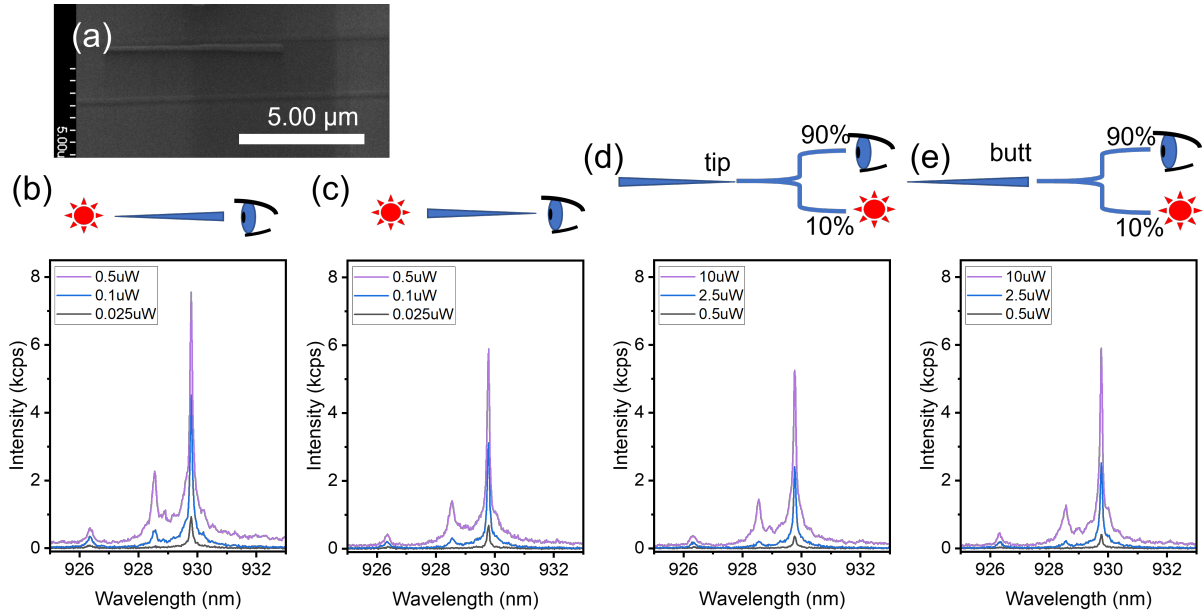


Figure 6.10: (a) SEM of nanowire on waveguide that has been broken in half. Pump and collections schemes for PL measurements on one waveguide. (b) Pumping from the base of the nanowire end and collecting from the tip. (c) Pumping into the tip and collecting from the base end. Using a 90:10 beamsplitter to pump and collect from (d) the tip and (e) from the base of the nanowire.

Before transfer, the QD exciton complex was emitting at 929.1 nm. After transfer, the QD was emitting at 929.8 nm. In each of the spectra, large phonon sidebands are observed even at lower pump powers. The large shift in wavelength and broadening of the emission line is indicative of insufficient cooling of the sample [15].

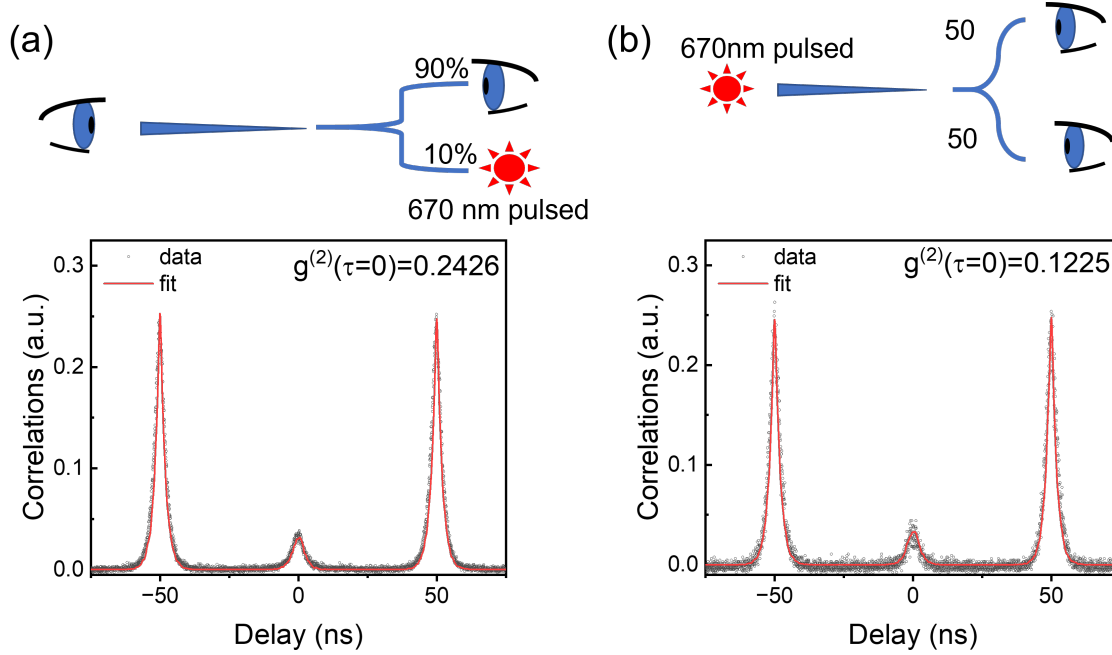


Figure 6.11: Autocorrelation ($g^{(2)}(\tau)$) measurement collecting from (a) both output ports of a device and (b) one output port and the transmission passing through a beam splitter and sent to two detectors.

6.3.3 Single Photon Purity

It can be seen that the power needed to saturate the QD is similar regardless of which end of the nanowire the excitation laser was pumped into. Using a 90/10 fibre power splitter to pump and collect on one side of the waveguide and another fibre to collect emission from the other end, the emission was sent to two tunable filters set to filter for the same emission complex at 930 nm and sent to two SNSPDs (see Fig. 6.11 (a)). A $g^{(2)}(\tau = 0) = 0.24$ was measured using this configuration.

We compare this $g^{(2)}$ with a conventional setup where the collected emission from one side is filtered for the desired complex and sent to a beamsplitter (see Fig. 6.11 (b)). This yielded a $g^{(2)}(\tau = 0) = 0.12$.

6.3.4 Sub-conclusion

Since the spectra looked much broader than pre-transfer, we realized the redshift and excess broadening were likely due to poor thermal contact with the system. The problem was not immediately obvious because the thermistor (ruthenium oxide) is located on the 4 K stage (circular gold-plated copper plate in Fig. 6.12) and the sample is mounted farther down on a gold-plated copper rod that extends several centimetres away from the plate. The temperature reading of 4 K stage was thus lower than the true sample temperature.

After opening the cryostat, we noticed the chip was poorly adhering to the gold-plated sample mount. It is worth noting that the fibre arrays remained securely attached to the

chip, even when the chip became completely detached from the mount and was temporarily supported solely by the fibres. The sample was remounted to the base using the an epoxy (3M Scotch-Weld epoxy adhesive 2216 B/A). Additionally, the excess bulk material from the gold-plated rod and copper block that the chip and holder were mounted to was making it difficult for the coldhead to sufficiently cool the sample. To mitigate this problem, a new, thinner, L-shaped sample mount was installed. The before and after versions of the sample stage are shown in Fig. 6.12. By remounting the sample and using a thinner sample stage, we were able to recover spectra that more closely matched the pre-transfer spectra.

In the following paper, we only present results using the new sample stage on the 4 K stage.

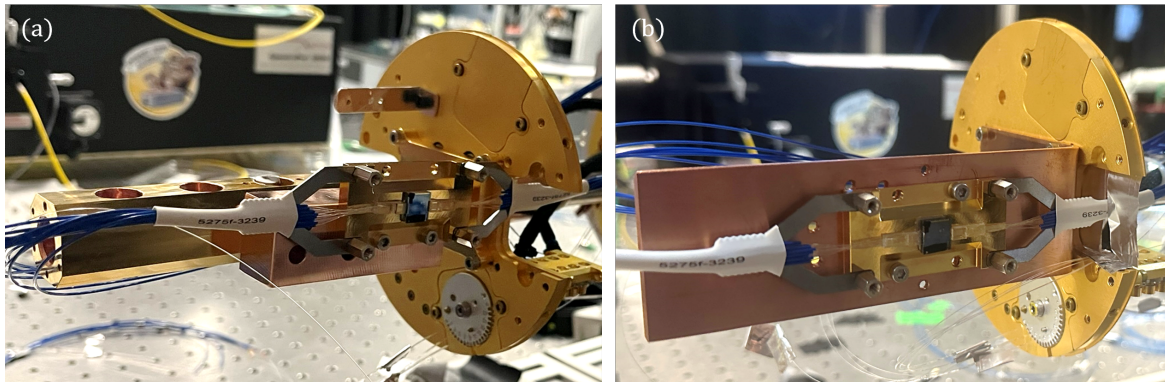


Figure 6.12: Photographs of the fibre-packaged chip on (a) the old sample mount connected to the 4 K stage featuring a long metal rod and copper block attachment and (b) the new sample mount that uses a thin L-shaped copper plate.

6.4 Article: Fibre-packaged plug-and-play single-photon sources using chip-integrated nanowire quantum dots

Adapted with permission from **Yeung, E.**, Zeeshan, M., Northeast, D. B., Weber, J. F., Poole, P. J., Williams, R. L., Dalacu, D. (2026). Plug-and-play single-photon sources using chip-integrated nanowire quantum dots. *Appl. Phys. Lett.* 129(8), 084005. <https://doi.org/10.1063/5.0348784>. Copyright 2026 AIP Publishing LLC.

6.4.1 Statement of Contributions

EY and DN ran simulations on COMSOL and Lumerical related to the chip design. EY performed the simulations necessary for determining the geometry of the waveguides. DN performed the additional simulations looking at the power transfer from the nanowire to the waveguide for different NW geometries and orientations. EY wrote the code and made

the layout for the chip submitted to LioniX International for fabrication. PHIX Photonics Assembly fibre-packaged the chip.

EY performed the loss measurements of the non fibre-coupled chip using lensed fibres whilst MZ performed the same set of measurements with cleaved fibres. JW, MZ, and DN spliced the fibres from the fibre-packaged chip to the fibre feedthroughs in the cryostat.

EY and MZ performed the optical measurements to characterize the devices.

EY, MZ, DN, RW, PP, and DD were in active discussions throughout the project. Data analysis and manuscript preparation was done by EY. DD supervised the research.

Plug-and-play Single Photon Sources using Chip-integrated Nanowire Quantum Dots

Edith Yeung,^{1,2} Mohd Zeeshan,² David B. Northeast,² John F. Weber,² Philip J. Poole,² Robin L. Williams,² and Dan Dalacu^{1,2}

¹*University of Ottawa, Ottawa, Ontario, Canada, K1N 6N5.*

²*National Research Council Canada, Ottawa, Ontario, Canada, K1A 0R6.*

We present a fixed-alignment platform for delivering single photons directly in fiber using chip-integrated nanowire quantum dots. Emission from III-V nanowire quantum dots is evanescently coupled to silicon-based integrated photonic circuitry. Fiber arrays fixed to the chip are used to efficiently and reliably couple light on- and off-chip allowing for plug-and-play operation of multiple sources on the same chip. We demonstrate plug-and-play in-fiber single photon collection efficiencies of $\eta_{\text{eff}} > 1.6\%$ from multiple devices on the same chip with multiphoton emission probabilities as low as $g^{(2)}(0) = 7.2\%$. Device performance was robust against multiple cool-down cycles and stable over tens of hours of continuous operation. Such plug-and-play fiber-coupled single photon sources will be a key resource required for distributing quantum information in future quantum networks.

Future photonic quantum technologies¹ in computation² and communication³ rely on the high-rate on-demand availability of indistinguishable single photons. Sources based on semiconductor quantum dots⁴ can generate such photons^{5,6} in device structures that can operate with high efficiency⁷⁻⁹. These current state-of-the-art sources are typically operated in low-vibration closed-cycle cryostats equipped with all components necessary for achieving active alignment of individual sources. For future widespread deployment, however, 'plug-and-play' operation will be required where state-of-the-art performance is reproduced in alignment-free fiber-coupled devices.

The potential for fiber-coupled plug-and-play operation was already demonstrated almost two decades ago^{10,11} and performance has significantly improved as more sophisticated alignment techniques were developed¹². Two general approaches have been pursued to address the high precision alignment required for efficient coupling to single mode fiber. One involves pick and place techniques whereby photonic structures containing quantum dots are picked up from the growth substrate and placed directly on the core of a cleaved fiber¹³⁻¹⁵ or on the taper of a pulled fiber¹⁶. In a second approach, fibers are fixed to semiconductor chips at room temperature using optical alignment techniques to position the fiber for optimal coupling to on-chip photonic structures¹⁷⁻²². Rapid progress has led to current state-of-the-art plug-and-play collection efficiencies in fiber of over 10%^{21,22}.

A natural next step to address scalability limitations of the above approaches would incorporate fiber arrays in place of the single fibers²³. For sources emitting out-of-plane, achieving optimal coupling for more than one device will be challenging due to the random nucleation position of the quantum dots. Sources emitting in-plane are better suited for coupling to fiber arrays with waveguide-fiber alignment facilitated through advanced fiber-integration methods such as photonic wire bonding²⁴. One would still have to rely on finding optimally positioned quantum dots somewhere along the

lengths of multiple adjacent waveguides. A more deterministic approach harnesses position-controlled nanowire growth²⁵ together with a pick and place technique to assemble hybrid devices where nanowire quantum dot sources²⁶ are integrated with Si₃N₄ waveguides²⁷⁻³⁰. Each waveguide is assured collection of the emission from a single quantum dot via efficient evanescent coupling between nanowire and waveguide modes.

In this work we report on a fiber-packaged silicon-based integrated photonic chip hosting multiple III-V nanowire quantum dot single photon emitters. We use commercially fiber-packaged chips and assemble the hybrid devices using a pick and place approach^{29,30}. The plug-and-play device performance was tested in a compact Gifford-McMahon cryo-cooler with a base temperature of 2.2 K. Throughput measurements with a laser revealed no significant fiber-waveguide misalignment during cool-down. Alignment was also stable over multiple cool-downs with only less than 4% variation in single photon count rate across 5 cycles and only 3.55% variation during 28 hours of continuous operation. Plug-and-play in-fiber single photon collection efficiencies of $\eta_{\text{eff}} > 1.6\%$ were obtained for multiple sources on the same chip.

Fabrication of the designed circuit was sourced from a commercial foundry (LioniX International) and included packaging of the fiber arrays for plug-and-play operation. A packaged chip with standard fiber arrays butt-coupled at each facet of the chip is shown in Fig. 1(a) where each array contains 8 fibers. The arrays are fixed at room temperature with alignment achieved using bespoke waveguide loop-backs on each side of the chip.

Low-pressure chemical-vapour-deposited Si₃N₄ was used for the photonic integrated circuit (PIC). The waveguides are 100 nm thick and 1 μm wide and are clad with 4 μm of SiO₂ above and 8 μm below. At the edge facets, the waveguide widths are tapered to 250 nm over a length of 200 μm for mode-matching to the single mode fibers. From simulations we obtain 88.6% (82.6%) TE (TM) mode overlap between the mode at the waveguide facet and PM780HP fiber mode with a 0.4% loss for both TE and TM modes as light propagates along the 200 μm

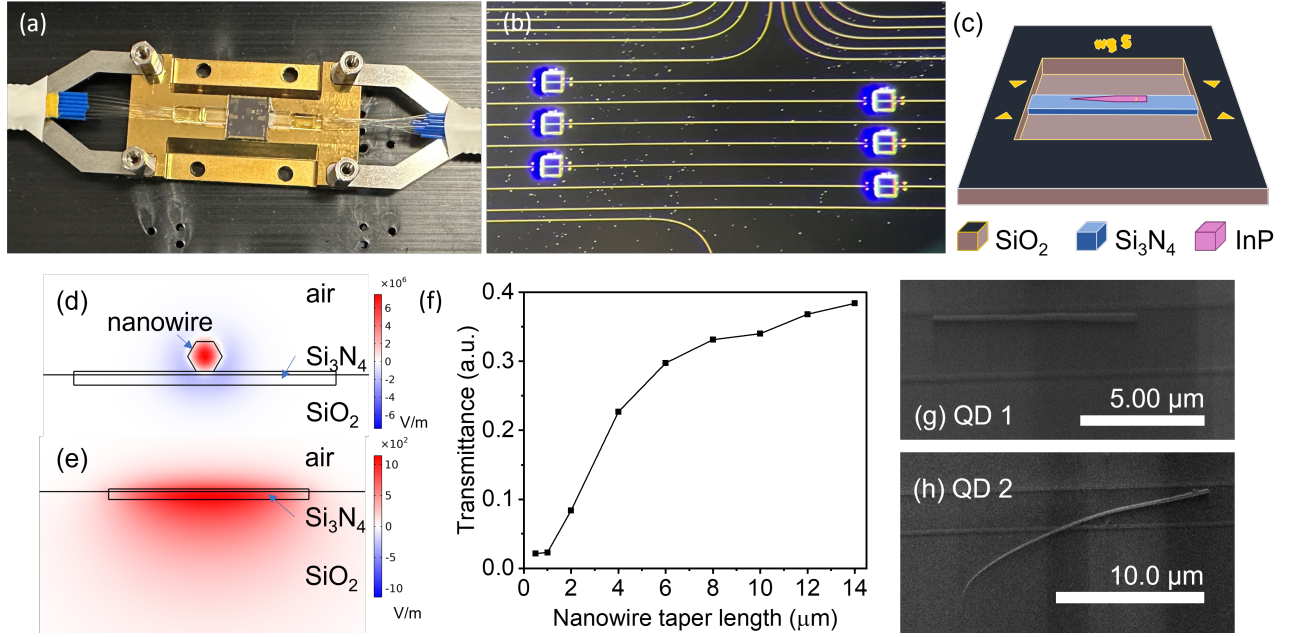


FIG. 1: (a) Optical image of the fiber-packaged chip. (b) Zoom in of the six device waveguides showing the windows for nanowire placement. (c) Schematic of the nanowire (purple) on the waveguide (blue) within the window. Gold markings were used to assist in nanowire placement. (d) Simulations showing the electric field through a cross-section of the Si₃N₄ waveguide in the window region with a nanowire on top. Only the top 25 nm of the 100 nm thick Si₃N₄ waveguide is exposed from the bottom SiO₂. (e) As in (d) but without the nanowire. (f) Calculated fractional power transferred from the nanowire mode in (d) into the waveguide mode in (e) as a function of the nanowire taper length. (g,h) SEM images of nanowires QD 1 and QD 2 transferred to waveguides.

length of the taper. Thus we expect an overall coupling efficiency of 88.2% (82.4%) for TE (TM) at the two facets.

For nanowire placement, windows were opened in the top oxide to expose the Si₃N₄ waveguide, shown schematically in Fig. 1(c) and as highlighted blue squares in Fig. 1(b). The oxide is etched to expose 25 nm of the waveguide, providing sufficient visibility to facilitate nanowire placement. As the waveguides approach the window regions they are tapered from a width of 1 μm to 2.2 μm, ensuring single TE-mode operation within the window region i.e. the TM mode is not confined. Using a taper length of 50 μm, we calculate a transmission loss of 0.1% for the TE mode. The simulated electric field of the TE mode confined in the Si₃N₄ in the window region is shown in Fig. 1(e).

Each nanowire to be placed hosts a single quantum dot that provides the single photons to the PIC. The nanowires are grown using position-controlled selective-area vapour-liquid-solid epitaxy described in detail in Refs. 25,26. The nominally ~3 nm thick InAsP quantum dot lies 1 μm from the base of a 20 nm diameter InP core, which is clad in an InP shell for a final base diameter of 250 nm. The structure tapers to a tip diameter of 20 nm along its ~20 μm length to expand the confined HE₁₁ mode and achieve coupling to the underlying Si₃N₄ waveguide²⁹.

To assemble the hybrid InP nanowire - Si₃N₄ waveguide device, we employ a pick-and-place technique described

in detail in Refs. 29,30. We first identify candidate nanowires using low temperature photoluminescence (PL) measurements in a confocal micro-PL (i.e. free-space) setup. These measurements are done in an ultra low-vibration closed-cycle helium cryostat equipped with high precision xyz piezo motors for alignment. The nanowires are selected based on brightness and spectra with a preference for dots emitting from just one or two charge complexes. Spectra of two selected nanowires measured on the growth substrate using the free-space set-up are shown in Figs. 2(a,b) below.

After free-space characterization, the nanowire chip is mounted in a scanning electron microscope (SEM)-based nanomanipulator together with the PIC chip. The nanomanipulator is equipped with tungsten tips mounted on piezo motors used to pick up individual nanowires from the growth substrate and place them on the Si₃N₄ waveguides on PIC chip. Transferring with real-time imaging in an electron microscope allows targeting nanowire placement to within a few tens of nanometers. Ideally, the nanowire is positioned parallel to the waveguide and at its centre, as shown in Fig. 1(d). For this ideal case, simulations (at $\lambda = 920$ nm) of the fractional power of the nanowire fundamental HE₁₁ mode coupled into the waveguide TE mode is ~38% for a nanowire with a sufficiently long taper, see Fig. 1(f). This calculated coupling is lower than previous simulations due to the non-ideal waveguide geometry, see Ref. 29.

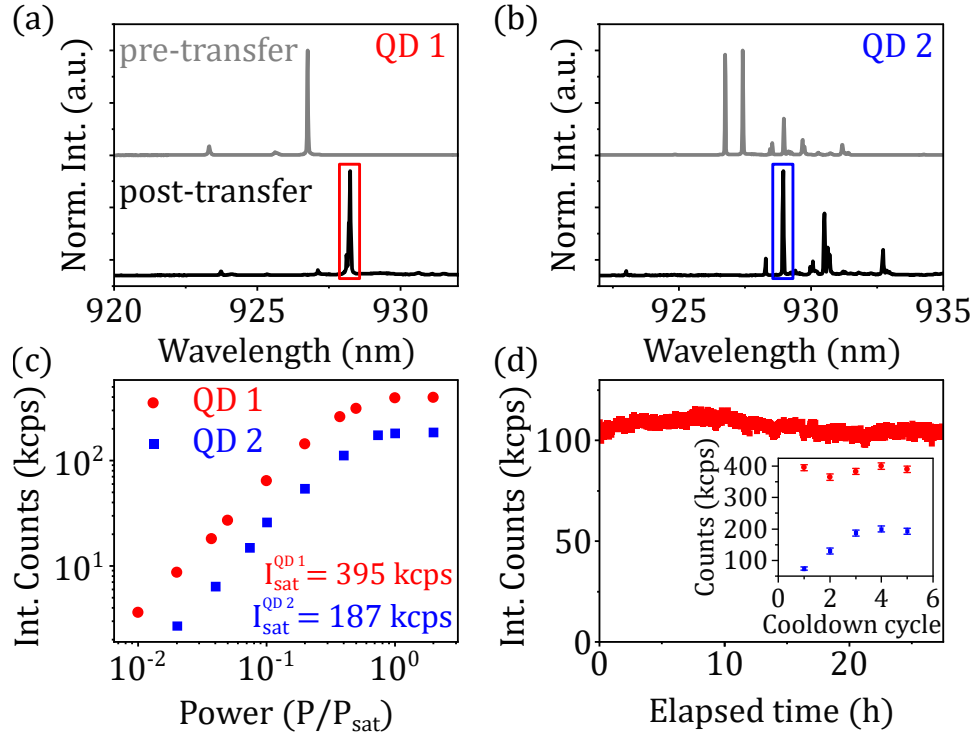


FIG. 2: (a, b) Emission spectra of two nanowire QDs measured before (grey) and after (black) transfer onto the Si_3N_4 waveguide fiber-packaged chip. In both cases there is a slight red-shift and redistribution of peak intensities. (c) Power-dependent PL peak intensities measured with the SNSPD on the third cool down. Counts at saturation are 395 kcps for QD 1 (red) and 187 kcps for QD 2 (blue) using 80 MHz above-band excitation. (d) Stability of the counts from QD 1 measured on the SNSPD over 28 hours. Inset shows count rates from QD 1 and QD 2 over multiple cool-down cycles spanning two months.

We attempted nanowire transfer to all six windows shown in Fig. 1(b) but due to the wider waveguides used here compared to our previous work³⁰, ideal nanowire placement proved challenging and we were only able to successfully transfer to half of the devices. We will focus on two of these, shown in Figs. 1(g) and (h), for which nanowire placement was not ideal but from which we still expect reasonable coupling. The first nanowire, QD 1 in Fig. 1(g), had the tapered tip broken off during the transfer process. Based on the simulation shown in Fig. 1(d), we estimate a drop in coupling efficiency to 34% due to the shortened, now $\sim 10 \mu\text{m}$, taper length³¹. More critical is the offset from the waveguide centre, which simulations predict a drop in the transmitted fractional power down to 26.9% for the observed 850 nm offset, see Supplementary Material.

The second nanowire, QD 2 in Fig. 1(h), is misaligned with respect to the waveguide such that the nanowire tip extends off the Si_3N_4 and onto the SiO_2 . There remains, however, a sufficiently long section, $\sim 9.1 \mu\text{m}$, of the nanowire on the waveguide from which we estimate a coupling efficiency of $\sim 34\%$ using again the simulations shown in Fig. 1(f). Adding the loss expected from the 12° misalignment across the top of the waveguide, the predicted coupling for this device is reduced to 23.7%,

see Supplementary Material.

For low temperature plug-and-play optical measurements, the packaged chip was mounted in a vacuum chamber cooled by a Gifford-McMahon cryo-cooler. This much simpler cryostat is only suitable for plug-and-play operation due to excessive vibrations of tens of microns with each refrigeration cycle. PM780HP fibers are fed into the cryostat via fiber feedthroughs and spliced to the fiber array of the packaged chip. The quantum dots were optically pumped with a 80 MHz pulsed laser at $\lambda = 780 \text{ nm}$ through one fiber with the emission collected from the second fiber on the opposite facet. Excitation was always via the base of nanowire i.e. from the right in Fig. 1(c), with collection via the tapered end i.e. using the fiber attached to the left facet.

PL spectra were measured with a fiber-coupled grating spectrometer equipped with a liquid nitrogen-cooled silicon charged-coupled-device. The PL spectra from QD 1 and QD 2 taken before (upper panel) and after (lower panel) transfer are shown in Figs. 2(a) and (b), respectively. In each case, sharp peaks are observed corresponding to emission from excitonic complexes in the QD. Post-transfer, the spectra of both dots are slightly red-shifted, by $\sim 1.5 \text{ nm}$, likely due to the modified strain environment after placement on the Si_3N_4 waveguide³². Also

evident are slight differences in the detailed structure of the spectra pre- and post-transfer. The single dominant peak at $\lambda = 927$ nm observed for QD 1 pre-transfer developed two smaller peaks on its high energy flank, seen more clearly in Fig. 3(a). QD 2 originally showed an equally weighted pair of dominant peaks at $\lambda = 927$ nm and $\lambda = 927.5$ nm whereas post-transfer, emission was primarily from the low energy peak. Such redistribution between different charge complexes suggest changes in the electrostatic environment surrounding the dot upon transfer to the Si_3N_4 waveguide. We note that similar redistribution was observed in Ref. 21 which the authors attributed to changes in the strain environment after attaching fibers.

To characterize the plug-and-play performance of the sources, we first filter the charge complex to be measured, highlighted in Figs. 2(a) and (b) by the coloured rectangles, using a fiber-coupled tuneable filter with a 0.1 nm bandwidth. The filtered emission is sent to either a superconducting nanowire single photon detector (SNSPD) or a Si avalanche photodiode detector (APD). In Fig. 2(c) we show the count rates measured on the SNSPD as a function of excitation power normalized to the power at which the intensity saturates, P_{sat} . For the highlighted peaks in Fig. 2(a,b), we measure count rates at P_{sat} of 395 kcps from QD 1 and 187 kcps from QD 2, where measurements were taken during the third cool-down, see below.

For determining the robustness of the plug-and-play device, the chip was cycled through five 300 K–2.2 K cool-downs. The inset of Fig. 2(d) shows the count rates for both quantum dots measured at saturation as a function of cool-down cycle. With the exception of the counts of QD 2 in first and second cool-down, we observe only small variations of 3.52% for QD 1 and 3.37% for QD 2 in count rates over multiple cycles. For the initial cool-downs, certain fibers required re-splicing due to accidental breakage, likely due to bad splicing, and this resulted in increased count rates for QD 2 after the second, and again after the third, cool-down.

On the fifth cool-down, we also evaluated the stability of the system by monitoring count rates over extended periods. Using a PID-controlled variable optical attenuator to maintain the laser input power at precisely $P_{\text{sat}}/8$, count rates from QD 1 were monitored over 28 hours. Only slight variations in counts were observed, see Fig. 2(d), and by fitting the histogram of the signal distribution with a Gaussian function, we obtained a standard deviation of $\sigma = 4.13$ kcps (3.55%), see Supplementary Material.

Next we estimate the collection efficiency obtained with plug-and-play operation, η_{eff} , defined here as the fraction of single photons per excitation pulse collected off-chip in fiber. For these measurements, we use APD count rates at saturation collected on the third cool-down: $I_{\text{meas}} = 210 \pm 5$ kcps for QD 1 and $I_{\text{meas}} = 129 \pm 5$ kcps for QD 2. Using a tunable laser set to around the emission of the QDs ($\lambda = 930$ nm), we measured

Loss source	QD 1	QD 2
Propagation loss	0.12 dB (97.3%)	0.12 dB (97.3%)
Facet loss	2.8 dB (52.5%)	2.8 dB (52.5%)
Single mode operation	3 dB (50%)	3 dB (50%)
Unidirectional operation	3 dB (50%)	3 dB (50%)
Nanowire misalignment	5.7 dB (26.9%)	6.25 dB (23.7%)
Filtered acoustic phonons	1 dB (80%)	1 dB (80%)
Total	15.62 dB (2.74%)	16.17 dB (2.42%)

TABLE I: Estimates of photon loss from different sources for the two devices. The propagation loss was calculated using $0.3 \text{ dB/cm} \times L$ where $L = 0.4$ cm is the distance from the QD to the facet edge when collecting from the tapered end of the nanowire.

a transmittance of the collection arm of $8.09 \pm 0.04\%$, where the losses are due to the tunable filter and associated mating sleeves. Assuming an APD efficiency of $25 \pm 5\%$ from the manufacturer, we estimate a corrected count rate out-coupled from the chip into fiber of $I_{\text{corr}} = I_{\text{meas}} / (0.0809 \cdot 0.25)$. The efficiency is calculated as $\eta_{\text{eff}} = I_{\text{corr}} / f_{\text{rep}}$ where $f_{\text{rep}} = 80$ MHz is the repetition rate of the pulsed laser. We obtain collection efficiencies of $\eta_{\text{eff}} = 2.7\%$ for QD 1 and $\eta_{\text{eff}} = 1.7\%$ for QD 2. These are reduced to $\eta_{\text{eff}} = 1.9\%$ for QD 1 and $\eta_{\text{eff}} = 1.6\%$ for QD 2 after correcting for multiphoton emission events, see below.

In Table I we list identified sources of photon loss with estimated values for each device. Propagation (scattering/absorption) and facet loss were obtained from room temperature measurements of cut-back structures on the chip, see Supplementary Material. The validity of using these room temperature values at 2.2 K was tested by monitoring the throughput of a selected waveguide during a cool-down cycle and observing negligible variation in output power, see Supplementary Material. Single mode operation results in the loss of photons generated with polarization orthogonal to the TE mode (i.e. there is no TM mode to couple to) and unidirectional operation results in the loss of photons emitted towards the base of the nanowire (i.e. in the direction opposite the collection fiber). The losses associated with reduced nanowire-waveguide coupling due to nanowire misalignment have been discussed above. The final 20% loss is from emission into the phonon sidebands³³ which have been filtered out by the 0.1 nm bandpass filter. Taking account of the loss listed, we expect efficiencies to be limited to 2.74% for QD 1 and 2.42% for QD 2, remarkably close to the measured values, $\eta_{\text{eff}} = 2.7\%$ and $\eta_{\text{eff}} = 1.7\%$, respectively.

These quoted single photon source efficiencies represent upper limits as we have not considered multiphoton emission probabilities. To verify the single-photon purity of the sources, we measure second-order correlations, $g^{(2)}(\tau)$, of each device using a Hanbury Brown & Twiss setup. The filtered emission is directed to two SNSPDs via a fiber-based 50:50 splitter and temporal correlations are recording with counting electronics. The pump power-dependent correlation measurements of the

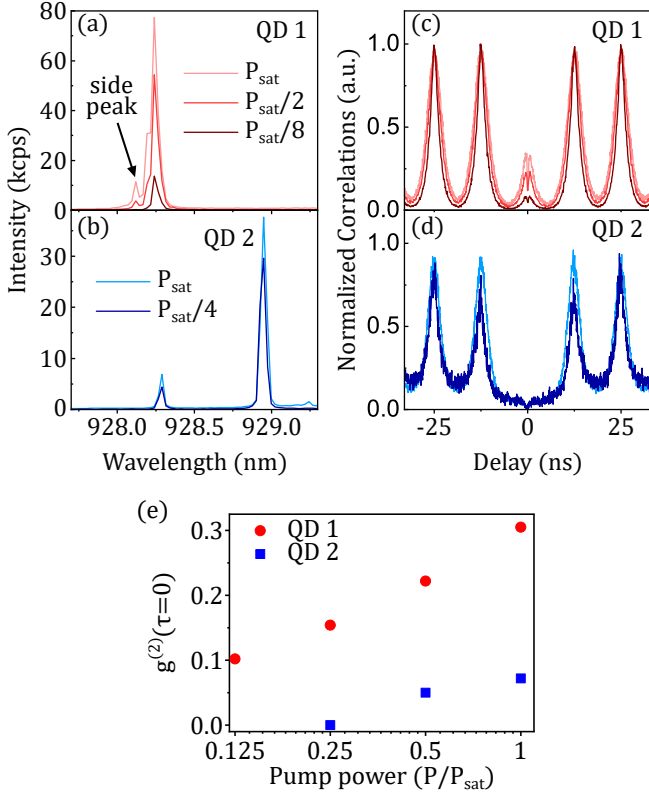


FIG. 3: Power-dependent PL spectra of (a) QD 1 and (b) QD 2. Power-dependent coincidences of (c) the $\lambda = 928.25$ nm peak from QD 1 and (d) the $\lambda = 929$ nm peak from QD 2. Both plots have been normalized to the area of the side peaks. (e) Extracted coincidences in the zero-delay peak, $g^{(2)}(\tau = 0)$, as a function of power.

filtered peaks of QD 1 and QD 2 are shown in Fig. 3(c) and (d), respectively.

For QD 2, the zero-delay peak is almost entirely absent, a signature of pure single photon emission. Coincidences at $\tau = 0$ drop to zero, rising slightly at small delays to show hints of the ‘volcano’ feature associated with re-excitation³⁴ typically present when exciting above-band. QD 1, on the other hand, shows significant coincidences in the zero-delay peak with a strong dependence on excitation power. Like QD 2, QD 1 also shows the $\tau = 0$ dip and associated side-lobes indicating re-excitation. The extracted power-dependent coincidences in the zero-delay peak normalized to the side-peaks, $g^{(2)}(\tau = 0)$, are plotted in Fig. 3(e) for both QD 1 and QD 2. Whereas $g^{(2)}(\tau = 0)$ rises linear from 0.1 to 0.3 as the excitation is increased from $P_{\text{sat}}/8$ to P_{sat} for QD 1, $g^{(2)}(\tau = 0)$ for QD 2 is negligible at low powers, rising only to 0.07 at P_{sat} .

To understand the different behaviour observed for the two devices, we consider the power-dependent PL spectra of both dots, shown in Figs. 3(a,b). In the case of QD 1, two additional peaks appear on the high energy flank of the original peak as the excitation power is increased, as mentioned previously. The tail of the closer

peak overlaps the original peak and cannot be filtered out with the tunable filter. The correspondence between the increasing intensity of this side peak and the increasing zero-delay coincidence counts as the excitation power is increased suggests that the two peaks are not mutually exclusive, i.e. the side peak pollutes the single-photon purity. Conversely, the measured peak in QD 2 remains a single, well-isolated peak for all excitation powers used, resulting in only a slight increase in zero-delay counts with power, likely entirely from re-excitation.

In summary, we have presented an alignment-free multi-emitter chip-based platform for delivering single photons in fiber. We achieved high performance plug-and-play operation using commercially fiber-packaged chips with negligible chip-fiber misalignment during cool-down cycles to cryogenic temperatures. The number of on-chip sources delivering single photons off-chip is limited only by available fiber-array sizes. Beyond single photon generation and distribution via fiber, the chip-based approach can be developed to include integrated photonic quantum computing³⁵ with outcomes available for distribution in future quantum networks.

Although in-fiber single photon collection efficiencies were limited to 1.6 – 1.9%, we expect future devices operating with efficiencies well in excess of 50% using modified PIC designs. Modifications include the use of waveguide architectures that collect the emission from both ends of the InP nanowire³⁶ with waveguides that support both TE and TM modes²⁹. Further, accurate nanowire-waveguide alignment can be achieved using lower aspect waveguide geometries with which we have demonstrated evanescent coupling efficiencies of $\sim 90\%$ ³⁰. These waveguide modifications alone provide a 12-15 $\times$ increase in the in-fiber collection efficiencies with further increases possible using better facet coupler designs.

See the Supplementary Material for additional simulations of InP nanowire-Si₃N₄ waveguide coupling, Si₃N₄ waveguide propagation and facet loss measurements, and temperature-dependent Si₃N₄ waveguide throughput measurements.

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AUTHOR DECLARATIONS

Conflict of Interest

The authors have no conflicts to disclose.

Author Contributions

Edith Yeung: Data curation (equal); Writing –

original draft (equal); Writing – review & editing (equal). **Mohd Zeeshan:** Data curation (equal); Writing – original draft (equal); Writing – review & editing (equal). **David B. Northeast:** Resources (equal); **John F. Weber:** Resources (equal); **Philip J. Poole:** Conceptualization (equal); Resources (equal); **Robin L. Williams:** Conceptualization (equal); Resources (equal); **Dan Dalacu:** Conceptualization (equal); Writing – review & editing (equal); Supervision (lead).

DATA AVAILABILITY

Data supporting the findings of this study are available from the corresponding author upon reasonable request.

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Plug-and-play Single Photon Sources using Chip-integrated Nanowire Quantum Dots

Edith Yeung,^{1,2} Mohd Zeeshan,² David B. Northeast,² John F.
Weber,² Philip J. Poole,² Robin L. Williams,² and Dan Dalacu^{1,2}

¹*University of Ottawa, Ottawa, Ontario, Canada, K1N 6N5.*

²*National Research Council Canada,
Ottawa, Ontario, Canada, K1A 0R6.*

I. NANOWIRE - WAVEGUIDE ALIGNMENT

Fig. S1 shows COMSOL simulations (performed at $\lambda = 920$ nm) of the electric field E_z at mid-waveguide height of the TE mode in the nanowire - waveguide device. The simulations were performed for different alignments of the nanowire with respect to the Si_3N_4 waveguide: (a) centred, (b) offset by 850 nm, and (c) at a 12° angle. These three cases represent the ideal nanowire position and actual positions of QD 1 and QD 2 nanowires, respectively. For the simulations, we used 100 nm thick and $2.2\ \mu\text{m}$ wide Si_3N_4 waveguides. For the nanowires, we assumed a base radius of $r_{\text{NW}} = 125$ nm and a length $l_{\text{taper}} = 10\ \mu\text{m}$ except in (c) where the taper is set to $l_{\text{taper}} = 12\ \mu\text{m}$, allowing for part of it to hang off the side of the waveguide while maintaining a $10\ \mu\text{m}$ length on the waveguide. We obtain nanowire to waveguide coupling efficiencies of 34%, 26.9% and 23.7% for (a) the centred, (b) the offset and (c), the angled nanowire, respectively.

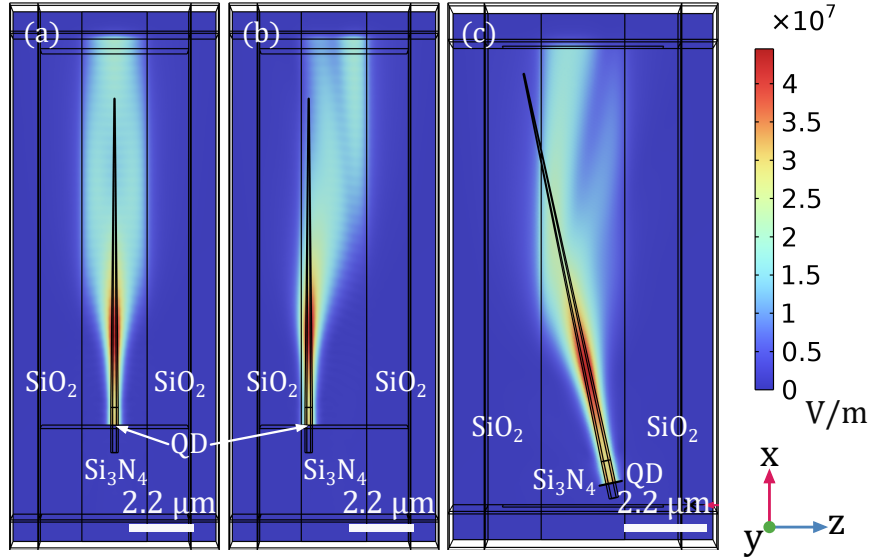


FIG. S1: Simulations of the TE mode in the nanowire-waveguide device for different nanowire alignment: (a) centred, (b) 850 nm offset, and (c) 12° angle, from which we calculate 34%, 26.9% and 23.7%, respectively, nanowire to waveguide coupling.

II. PROPAGATION AND COUPLING LOSS

Propagation and facet loss in the Si_3N_4 waveguides was extracted from transmission measurements through waveguides of different length (i.e. cut-back measurement). The

measurements were made for wavelengths 910 nm to 980 nm using a tunable CW laser polarized to couple either to the TE or the TM waveguide mode. We measured both a bare chip (i.e. not fiber-packaged) with and without windows as well as the fiber-packaged chip. For the former, lensed-fibers were used to couple on- and off-chip.

For the bare chips, three waveguides without window openings were measured with lengths 100,445 μm , 50,672 μm , and 20,442 μm . Six waveguides with window openings were measured: four repeats with length 9,682 μm and one repeat each of lengths 50,072 μm and 100,234 μm . A calibrated (i.e. system loss subtracted) measurement for the set of waveguides without windows is shown in Fig. S2(a) for the TE mode. To extract the wavelength-dependent loss, the transmitted power spectrum is fit with a fifth-order polynomial. The length-dependent loss for three wavelengths are plotted in Fig. S2(b) with the slope from a linear fit giving the propagation loss whilst the loss/facet is obtained from the y-intercept divided by two. In Fig. S2(c), we show the results for measurements on the waveguides with windows. The extracted propagation and facet losses for each set of waveguides as a function of the input wavelength is summarized in Fig. S2(d). We observe similar propagation loss for waveguides with and without windows, within $\sim 0.1 \text{ dB/cm}$, and $\sim 1.7 \text{ dB}$ larger facet loss for waveguides with windows, likely from scattering at the window edge and not associated with the actual edge facet of the chip.

In Fig. S3(a) we show the results of the same transmission measurements performed on the fibre-packaged chip before and the chip was inserted into the SEM for nanowire transfer. Unlike the lensed-fiber measurements above, here we observe significant wavelength-dependent oscillations, likely Fabry-Pérot fringes between the cleaved fiber facet and the chip facet. We estimate the wavelength-dependent transmission by fitting a fifth order polynomial to the data (shown in dark hue) intended to average out the oscillating values.

From the fits to the data in Fig. S3(a) we obtained the wavelength-dependent transmittance from each waveguide and used a linear fit to this data to extract the propagation and facet losses, as above. The results are shown in Fig. S3(c). The propagation loss between the two sets (i.e. lensed-fibers versus fixed cleaved fiber) is similar, $\sim 0.3 \text{ dB/cm}$, as expected since the propagation loss should not depend on the coupling method. We do observe some unexpected variation in the propagation loss at wavelengths $\lambda > 950 \text{ nm}$ which may be associated with an uncertainty in transmission measurement due to the Fabry-Pérot fringes. As expected, the facet loss in the packaged chip are significantly reduced, by over

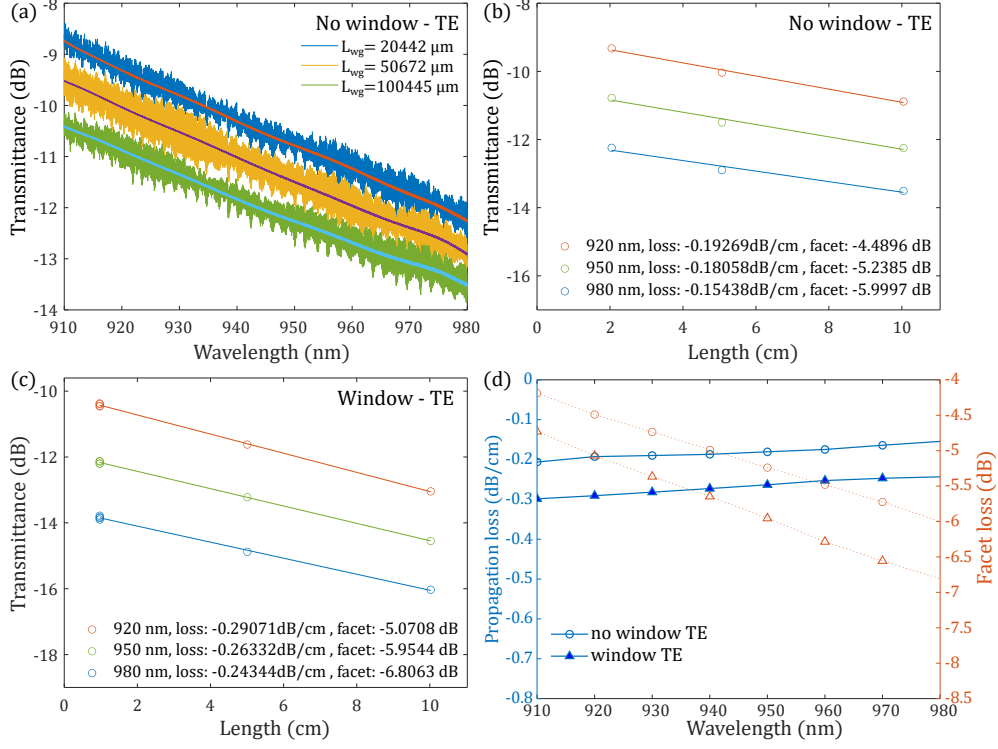


FIG. S2: (a) TE-polarized transmission measurement using lensed-fibers. Solid lines are fits to the data using a fifth-order polynomial. (b) Extracted wavelength-dependent loss as a function of waveguide length for waveguides (b) without and (c) with windows. Legends show propagation loss and loss/facet extracted linear fits to the data. (d) Extracted propagation loss (blue) and facet loss (orange) as a function of the wavelength.

-2 dB/facet: the couplers were designed for mode matching to cleaved PM780HP fibers. We use the propagation and facet loss from Fig. S3(c) to estimate collection efficiencies given in Table 1 of the main text.

To verify that the pick and place process in the SEM does not effect the performance of the fibre-coupled chip, we repeated the loss measurements post-transfer. These are shown in Fig. S3(b). The two longer waveguides, $L_{wg} = 100,234 \mu m$ and $L_{wg} = 50,072 \mu m$, which did not have nanowires transferred to them, showed very similar transmission measurements pre- and post-transfer. Both of the shorter waveguides, $L_{wg} = 9,062 \mu m$, did have nanowires transferred, and one of these (wav 5, QD 1) showed a clear reduction in transmission (~ 2 dB) post-transfer, expected from additional scattering due to the nanowire.

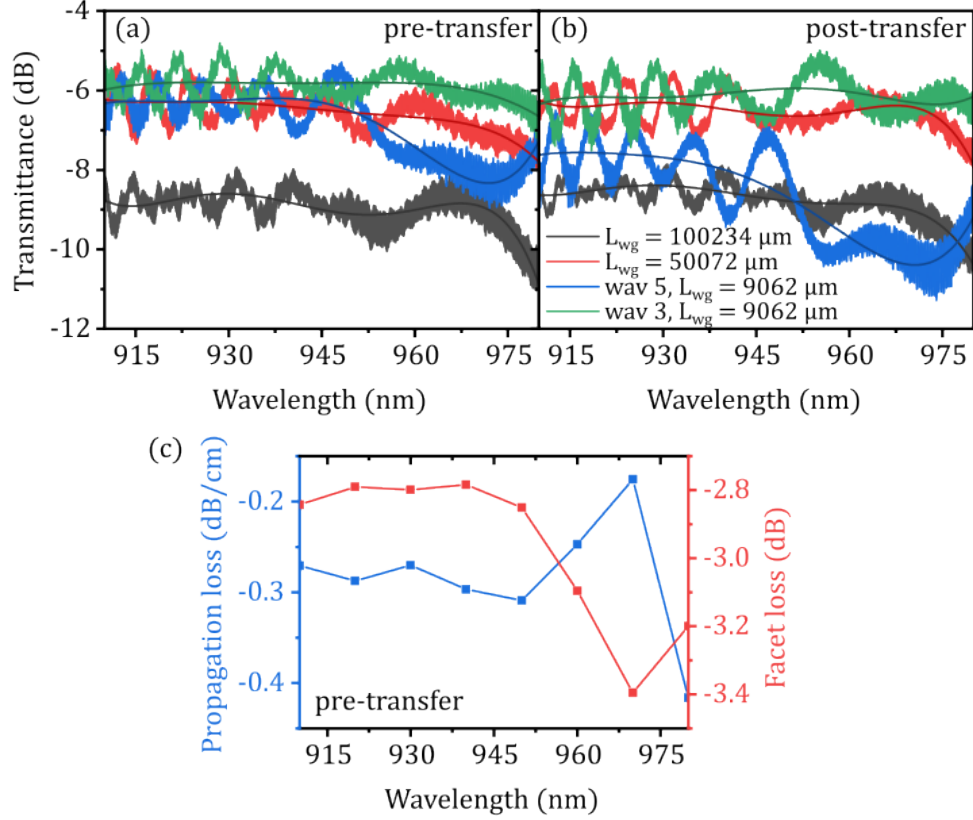


FIG. S3: Room-temperature TE-polarized transmission measurement of the fibre-packaged chip for waveguides of different lengths both (a) pre- and (b) post-nanowire transfer. Fits using fifth-order polynomials, shown in darker hue, are used as a rough estimate of the wavelength-dependent transmission. (c) Wavelength-dependent propagation loss (blue) and loss/facet (red) calculated using the fit to the data in (a).

III. TEMPERATURE-DEPENDENT THROUGHPUT

To verify the temperature stability of the fiber-waveguide coupling in the fixed-fiber device, we monitored the throughput of a selected waveguide during one of the cool-down cycles. The laser power measured at the output fiber for the TE mode launched into the waveguide via the input fiber is shown in Fig. S4. We observe temperature-dependent oscillations of $\sim 30\%$ in the measured power, likely the same Fabry-Pérot fringes observed above but here due to contraction of the cleaved fiber-chip facet spacer. Upon reaching the base temperature of 2.2 K, we measure, within a few percent, the same laser power as observed at 300 K.

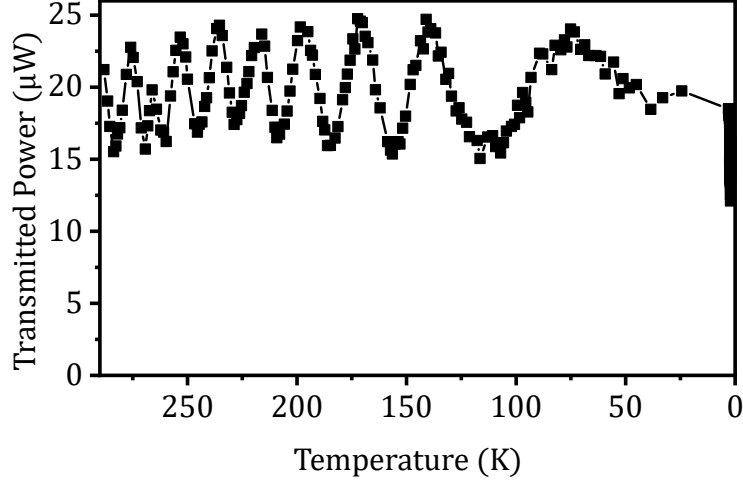


FIG. S4: Temperature-dependent transmission through a waveguide of the fiber-packaged chip. The measurement was made using a TE-polarized laser at $\lambda = 930$ nm.

IV. STABILITY OVER 28 HOURS

To determine the stability of the source, we monitored the intensity of the selected emission peak from QD 1 over 28 hours. To quantify the stability, we normalize the collected photon counts over 28 hours by the mean (107.3 kcps at $P = P_{\text{sat}}/8$) and bin the resulting normalized data to find the standard deviation $\sigma = 3.55\%$, see Fig. S5.

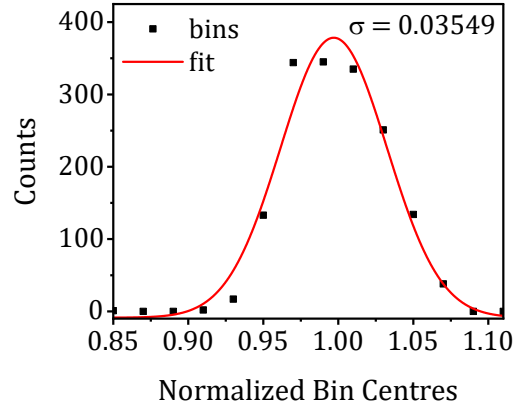


FIG. S5: Normalized counts relative to the mean for a 28 hour measurement of QD 1 (Fig. 2(d) of main text) giving a standard deviation of 3.55% .

6.5 Conclusion

We have demonstrated a plug-and-play single-photon source robust against multiple cooldown cycles. By eliminating the need for active alignment, our approach opens up new possibilities for a miniaturized, scalable plug-and-play platform for multiple emitters on chip.

Future work with waveguide geometries more suited for light with wavelengths around 950 nm would further improve the collection efficiency at first lens as shown in Chapter 4. It would be preferred to design PICs able to maintain both TE and TM modes which would increase the collection efficiency of both NW-QD HE_{11} modes and thus achieve a higher PnP efficiency. The waveguides enabling these improvements would be made from thicker Si_3N_4 films, enabling narrower widths within the placement window areas. These narrower geometries would also assist in the accurate placement of nanowires to locations providing better mode overlap to waveguide modes.

Although the long tapered geometry ($\sim 20 \mu m$) of the nanowire is more suited for free-space coupling to an objective, this geometry was not ideal for applications requiring the mechanical manipulation of the nanowire. This is due to the thin nanowire tip's propensity to adhere strongly to the SiO_2 surface through electrostatic forces and van der Waal forces. The smaller diameter made the nanowire more flexible when compared to nanowires with shorter tapers, which tend to be more rigid. Rather than roll across the oxide surface, these longer and thinner nanowires would bend and often stick to the transfer window surface which made transfer to a precise location more difficult.

Future iterations of the work will use nanowires with geometries more suited for the transfer process whilst maintaining sufficient power transfer from the NW to the waveguide.

A key aspect of this project is to have a single-photon source that emits in-plane into PIC waveguides. This allows us to use integrated technologies to build more complex photonic circuits for different applications. This work demonstrates a robust platform for stable light coupling on and off chip using NRC-grown NW-QDs.

Chapter 7

Conclusion

Over the course of my Ph.D., I have made significant progress in the hybrid integration of III-V quantum-dot nanowires and Si_3N_4 waveguides. Prior to this Ph.D., in 2016, preliminary work for a hybrid structure from Elshaari *et. al.* [110] and Zadeh *et. al.* [29] used a nanomanipulator to pre-position the nanowires on a silicon substrate then encase them in a Si_3N_4 waveguide. The collection efficiency from these early devices were $\sim 24\%$. Later in 2020, Mnaymneh *et. al.* [11] were the first to utilize an evanescent coupling approach with the nanowires by placing them outside pre-fabricated waveguides and achieving a coupling efficiency of 74%. These devices saw a 22 nm shift in the emission wavelength due to the encapsulating SiO_2 layer deposited after transfer.

Since then, we have moved on to using Si_3N_4 waveguides from commercial foundries and incorporating the NW-QDs adjacent to the waveguides through our pick-and-place method developed during my M.Sc. and mastered during my Ph.D. Utilizing this pick-and-place technique, I demonstrated successful integration with a wider variety of waveguide structures for other applications. Hybrid devices that I assembled have since been used by collaborators for applications in quantum memory [111]. The culmination of this work produced a true plug-and-play single-photon source.

7.1 Achieved results

In the first project presented in Chapter 4, we utilized a simple straight waveguide and tapered nanowire configuration to test the optical properties of a single-photon source device. Our experimental findings were in agreement with simulated results showing strong coupling efficiency dependent on the length of the nanowire taper. We were also able to show an improvement in the single-photon purity, lifetime, and two-photon indistinguishability using quasi-resonant excitation.

In Chapter 5, we moved on from using a tapered nanowire which restricted efficient collection from one end of the nanowire to an untapered nanowire with the QD positioned midway along the length of the nanowire adjacent to a curved waveguide. This configuration maintained the basic principles of evanescent coupling as before. Being a waveguide

itself, a NW-QD was used to bridge the gap between two Si_3N_4 waveguides. We performed bi-directional collection of the QD emission using this configuration and demonstrated that the optical properties of the QD were intact even with this geometry.

The goal of having a plug-and-play single-photon source using the NRC grown NW-QDs was finally realized in Chapter 6. In this project, I contributed to every stage of the project from planning (simulations and chip layout) to designing the experiments and measuring the devices. In the end, we were able to achieve a complete packaged product that could generate single-photons on-demand in a plug-and-play manner.

The scope of this PhD was very ambitious. There are still many more Si_3N_4 waveguide structures that I have tried to incorporate the nanowires onto that were not presented in this thesis. Among them included the incorporation of the nanowires into on-chip MZIs (both with and without gaps) and ring resonators. Unfortunately for the ring-resonator experiments, without the ability to controllably tune either QD emission or the resonances, the ring would act more as a filter than a means for Purcell enhancement. The lifetimes post-transfer demonstrated longer lifetime, indicative of a Purcell inhibition.

7.2 Future research directions

Our goal is to advance toward integrating telecom-band quantum dots (operating in the O and C bands) onto photonic integrated circuits to ensure compatibility with existing telecommunication infrastructure. Although this thesis does not present those results, preliminary studies on telecom devices were conducted in the early stages of my Ph.D. Since the initial versions of the nanowire quantum dots exhibited relatively low emission rates, my research focused on 950 nm nanowire quantum dots, which consistently demonstrated cleaner emission lines and improved photon collection [31, 15, 16]. Because of equipment failures that limited the growth of additional high-quality telecom nanowire quantum dots and the long turnaround time between design submission and receiving photonic integrated circuits from the foundries, there was insufficient time to pursue this direction further.

A logical next step would be to use waveguides fabricated in-house through a collaboration between the NRC and Carleton University. This approach would significantly shorten the turnaround time from design submission to final chip fabrication. During my Ph.D., I had the opportunity to contribute to this initiative by testing some of the fabricated chips. However, at that time, waveguides with window openings were not yet available. I was also briefly involved in the U-groove project, which aimed to develop an NRC-made plug-and-play source. In this work, cleaved fibres were successfully glued into U-grooves, and the transmission through the waveguide was monitored during a cooldown cycle. The preliminary results were encouraging, as the fibres remained firmly bonded to the chip throughout the process. It is worth noting that these initial chips did not include windows for nanowire placement.

With the capability to fabricate high-quality waveguide devices in-house, it will become possible to optimize the design of both the waveguides and nanowires for improved compatibility. As was mentioned for the plug-and-play source, the waveguides were a lot

wider and thinner than the waveguides presented in the earlier chapters. Due to the design limitations, we were unable to achieve higher collection efficiencies. By developing waveguides with integrated U-grooves and windows internally, we can greatly accelerate the fabrication process and realize devices with significantly higher end-to-end collection efficiencies.

The long-term goal for quantum integrated photonic circuits is to create a fully functional laboratory on a single chip that can generate, manipulate, and detect single photons within one integrated platform. This thesis specifically focused on the on-chip generation of single photons. There is a vast library of photonic structures that have been well studied including ring resonators (add-drop filters), beam-splitters (directional couplers), and phase shifters. During my Ph.D., I have assembled several devices that uses these more complex structures. However, without the ability to tune the emission wavelength of the quantum dots, these studies provided limited insight. In addition, the experimental setup allowed only a single lensed fibre on each side of the chip, which made the directional coupler design largely impractical. Ongoing research is being conducted to grow SNSPDs on PICs. A combination of the single-photon source, filter, directional couplers, and detectors on-chip would allow us to perform a full $g^{(2)}(\tau)$ measurement.

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