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Nonlinear Single-photon Generation for Photonic Quantum Technology

A thesis submitted for the degree of Doctor of Philosophy

by

Matthew Collins

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Australia

2015
To my family

Renee, Mum and Dad

Your endless love and support make anything possible!
Declaration of originality

To the best of my knowledge, this thesis contains no copy or paraphrase of work published by another person, except where duly acknowledged in the text. This thesis contains no material that has been previously presented for a degree at the University of Sydney or any other university.

I declare this thesis to be wholly my own work, unless stated otherwise.

Matthew J Collins
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Abstract

Single photons are the smallest indivisible quanta of light, canonically described by quantum mechanics. By carefully controlling the interaction of single photons, exquisite non-classical phenomena can be observed. Mature photonic chip technology has recently emerged as an ideal platform for quantum information processing using single photons. However, generating single photons efficiently on-chip remains a fundamental challenge. One solution is to harness the intrinsic nonlinearity available in certain photonic materials for nonlinear photon generation directly in on-chip waveguides themselves.

This work examines nonlinear photon generation in two key material platforms. The first is chalcogenide glass. Chalcogenide, while highly nonlinear, is amorphous and thus has broadband Raman noise. In this study the Raman noise is characterised at the single-photon level to find an intrinsic minima, which is then targeted for low-noise photon generation using an engineered waveguide. The second platform is silicon. As silicon is complementary metal-oxide-semiconductor (CMOS) fabrication compatible, it is congruent with mass production. Thus, in this study, photon-pair generation is first shown in a compact photonic crystal, before combining two monolithic sources using active multiplexing.

This thesis presents significant progress towards a key goal of the field – on-demand photon generation in a fully integrated photonic quantum processor.

![Artist’s impression of nonlinear photon-pair generation as a light pulse passes through a photonic crystal waveguide. Image courtesy of Dr. Joel Carpenter.](image-url)
Publications included in this thesis


List of abbreviations

1D one-dimensional
2D two-dimensional
ASE amplified spontaneous emission
ASMUX asymmetric spatial multiplexing
ATT attenuator
AWG arrayed waveguide grating
BPF band-pass filter
BS beam splitter
CAR coincidence to accidental ratio
CMOS complementary metal-oxide-semiconductor
CROW coupled resonator optical waveguide
CUDOS Centre for Ultrahigh bandwidth Devices for Optical Systems
CW continuous wave
CZ controlled phase
DG delay generator
DRW direct-write waveguide
EDL electronic delay line
EM electromagnetic
FBG fibre Bragg grating
FCA free-carrier absorption
FCD free-carrier dispersion
HBT Hanbury Brown and Twiss
HOM Hong-Ou-Mandel
KLM Knill-Laflamme-Milburn
LCoS liquid-crystal-on-silicon
LOQC linear optical quantum computing
MIR mid-infrared
MZI Mach-Zehnder interferometer
NA numerical aperture
ODL optical delay line
PCF photonic crystal fibre
PhCWG photonic crystal waveguide
PLZT lead lanthanum zirconium titanate
PPLN periodically poled lithium niobate
QIP quantum information processing
QC quantum cryptography
SEM scanning electron microscope
SFWM spontaneous four-wave mixing
SMUX symmetric spatial multiplexing
SOI silicon-on-insulator
SPD single-photon detector
SPDC spontaneous parametric down conversion
SpRS spontaneous Raman scattering
SSPD superconducting single-photon detector
TBPF tuneable band-pass filter
TE transverse electric
TIA time interval analyser
TM transverse magnetic
TPA two-photon absorption
WDM wavelength division multiplexer
XPM cross-phase modulation
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Chapter 1

Introduction

‘Nothing is impossible. Not if you can imagine it. That’s what being a scientist is all about.’

– Professor Hubert Farnsworth
Quantum photonics is an exciting field that combines fundamental science and discovery with real-world technological applications. Quantum particles have strange properties that do not have any analogue in classical physics, and do not behave like anything that people encounter in their daily lives, thereby leaving little basis for intuition. For example, if a quantum particle can exist in two possible states, such as up and down, is it possible that it can be both up and down at the same time? This is a quantum superposition. If quantum particles are used to store and process information, they are not limited by the standard rules of classical computation. It was shown in the early 1990s, that to solve certain classes of problems a ‘quantum computer’ can greatly out-perform a classical computer.

Any type of quantum particle can be used for a quantum computer, such as an electron, atom or photon. Photons do not interact strongly with matter, or with each other, a property ideal for storing and transmitting quantum information. This advantage of lack of interaction is why photons remain a prime candidate for the base particle of a quantum computer. Seminal work was done using table-top optical elements in the late 1990s and early 2000s. However, the field slowly realised that this kind of approach would become too cumbersome and unstable to scale up. Thus, quantum photonics emerged as a solution to the challenges of conventional quantum optics experiments.

Recently, it was shown that integrating an experiment onto a single millimetre scale monolithic chip provided unprecedented performance. Furthermore, the maturity of photonic fabrication technology offered the scalability necessary for complex circuits. Thus, the challenge then became to try to generate single photons on-demand in the photonic chip itself. Photons generated on-chip would have minimal loss between generation, processing and detection, thereby creating very efficient devices.

In nonlinear material, a strong coherent laser pulse can be used to generate correlated photon pairs. If one photon from this pair is detected, it heralds the existence and timing of the remaining photon; this is called ‘heralded single-photon generation’. The same strong coherent laser light generates uncorrelated photons, at frequency shifts corresponding to the characteristic energies of the molecular vibrations in the material, called ‘spontaneous Raman scat-
tering’ (SpRS). In amorphous materials, the SpRS spectrum is broadband and can add noise to the heralded single photon spectrum.

This thesis contributes to the field of quantum photonics by investigating two key questions in nonlinear photon generation. The first question asks: Is chalcogenide glass a viable platform for on-chip nonlinear photon generation? Chalcogenide (specifically As$_2$S$_3$ in this study) is a highly nonlinear glass and mature photonic integration platform that faces the challenge of uncorrelated SpRS noise. In Chapter 3, cooling is shown to provide limited reduction in the SpRS, measured indirectly through photon-pair correlations. To directly observe the SpRS and the effect of cooling, a novel photon-counting Raman spectrometer is designed and used to measure the Raman spectrum. An intrinsic minimum in the chalcogenide Raman spectrum is identified, called the ‘low-Raman window’, and targeted for low-noise photon generation in an engineered chalcogenide waveguide.

The second question asks: Can multiple nonlinear photon sources be combined using active spatial multiplexing to improve the overall photon statistics? Chapter 4 begins with nonlinear photon generation in compact silicon photonic crystal waveguides. Two monolithic devices are then combined using active spatial multiplexing. The output is shown to have significantly improved photon statistics compared to a single source.

In summary, chalcogenide is shown to be viable for on-chip nonlinear photon generation only in the limited region of the low-Raman window. In addition, active multiplexing clearly improves the statistics of photon generation by deterministically combining the output of two sources. Both of these investigations provide useful knowledge for the field to expand in order to work towards the goal of on-demand photon generation for photonic quantum processors.
Chapter 2

Photons and photonics

This chapter discusses the physical concepts and prior literature required to understand and contextualise the original work presented in this thesis.

‘People assume that time is a strict progression of cause to effect, but *actually* from a nonlinear, non-subjective viewpoint – it’s more like a big ball of wibbly wobbly... timey wimey... stuff.’

– Dr. Who
2.1 Quantum nature of light

In 1905 Albert Einstein first proposed the idea that light is in fact quantised into indivisible packets that we now call photons [1]. It took until the 1970’s before irrefutable evidence of the quantised nature of light was observed [2]. The photon is now recognised as a powerful tool for quantum information processing, with many proposals for quantum computing and communication [3]. During the last decade it has become clear that, to continue increasing the complexity of experiments and satisfy stability requirements, the field must move towards integrated photonics [4,5].

This section begins by briefly retracing the journey from the first experimental hints of the quantised nature of light to modern on-chip quantum information processing using single photons.

2.1.1 The breakdown of classical optics

Classically, light can be described as an electro-magnetic (EM) wave. This can be used to predict macroscopic behaviour, such as the diffraction of a finite wave as it travels through free space, or the interference of two waves, first shown by Young and the double slit experiment. In the late nineteenth century, James Clerk Maxwell established a set of equations that offered a complete description of the classical behaviour of EM waves – light. The Maxwell equations have the form

\[ \nabla \cdot \tilde{D} = \rho, \]

\[ \nabla \cdot \tilde{B} = 0, \]

\[ \nabla \times \tilde{H} - \frac{\partial \tilde{D}}{\partial t} = \tilde{J}, \]

\[ \nabla \times \tilde{E} + \frac{\partial \tilde{B}}{\partial t} = 0, \]

where \( \tilde{E} \) is the electric field, \( \tilde{H} \) is the magnetic field, \( \tilde{D} \) is the electric displacement field and \( \tilde{B} \) is the magnetic induction field. \( \tilde{J} \) and \( \rho \) are the free current and charge densities, respectively [6]. These equations remain invaluable for describing the majority of optical phenomena under investigation today.
At the end of the 19th century, classical optics failed to correctly describe the emission spectrum of a blackbody, a prediction known as the ultra-violet catastrophe. The solution was put forward by Max Planck, a German scientist, who at that time was a professor at the University of Berlin. He proposed that the energy of the oscillating electrons bound at the wall of the blackbody was quantised and that the energy in each frequency component of the radiated spectrum was related to this quantisation \[7,8\]. Prior to Planck, the Rayleigh-Jeans Law was the best attempt to describe the blackbody radiation emitted from a cavity. Classically the number of standing wave modes inside a cavity, per unit length and per unit frequency \(d\nu\), is given by \(N = \frac{8\pi\nu^2}{2c^3}\). At that time, the theorem of equipartition of energy stated that each mode of a cavity in thermal equilibrium has equal energy, \(E = k_BT\), where \(k_B\) is the Boltzmann constant and \(T\) is the temperature of the cavity/black-body. The Rayleigh-Jeans Law thus gave the energy density of radiation \(u(\nu)\), from the cavity, with

\[
u(\nu) = \frac{8\pi\nu^2}{c^3} k_BT, \tag{2.5}\]

where \(c\) is the speed of light. This equation predicted an infinite amount of energy is radiated from the cavity shown in the red curve in Fig. 2.1a). Planck’s proposal implied that the occupation of each mode follows the Bose-Einstein population distribution, 

\[
N_{BE} = \frac{1}{(e^{\frac{h\nu}{k_BT}} - 1)},
\]

with each mode having discrete energy \(E = h\nu\). The Planck theorem was thus

\[
u(\nu) = \frac{8\pi\nu^2}{c^3} \frac{h\nu}{e^{\frac{h\nu}{k_BT}} - 1}. \tag{2.6}\]

The blackbody emission spectrum predicted by Planck, seen in the blue curve in Fig. 2.1c), agrees with the Rayleigh-Jeans curve at low frequencies. However, it is qualitatively very different for the rest of the spectrum, as seen in Fig. 2.1b). The Planck equation matched the experimental observations of the black body emission from hot objects, such as stars. This discovery was central to the catalyse the study of quantum mechanics. However, both Planck’s work and Einstein’s photoelectric effect can be described semi-classically. In semi-classical physics, matter states are quantised, while light is described classically. While these experiments showed that something must be quantised,
there was not yet sufficient evidence to indicate that it had to be the light.

Optical coherence

In 1956, astronomers Robert Hanbury Brown and Richard Twiss devised a new type of stellar intensity interferometer that solved the emerging problem of how to stabilise collection mirrors across very long base lines [9]. For laboratory investigations the Hanbury Brown and Twiss (HBT) interferometer can be constructed by shining a light source onto a half silvered mirror, also known as a 50:50 beam splitter (BS), with one half of the beam incident onto a photo-detector and the other beam incident on a second photo-detector. This experiment showed that light originating from a star was phase correlated, or coherent, even though it was generated from random emissions of an atomic bath.

In 1963 Roy Glauber, a Harvard Professor of physics, published the optical theory of coherence and was co-awarded the 2005 Nobel Prize in Physics for his works. The first-order correlation function, $g^{(1)}(\tau)$, refers to the phase coherence of light, and quantitatively the contrast (known as visibility) in the interference pattern of a Michelson interferometer. The $g^{(1)}(\tau)$ can be expressed with the normalised field form [10]

$$g^{(1)}(\tau) = \frac{\langle E(t)\,^*\,E(t+\tau) \rangle}{\langle E(t)\,^*\,E(t) \rangle}.$$  \hspace{1cm} (2.7)
A truly monochromatic light source will have an infinite coherence time and a $|g^{(1)}(\tau)| = 1$. A real world light source with a frequency bandwidth of $\Delta \omega$ will have a coherence length of $1/\Delta \omega$ and $0 < |g^{(1)}(\tau)| < 1$. Incoherent light has a zero coherence length and $|g^{(1)}(\tau)| = 0$ [11].

The second-order correlation-function $g^{(2)}(\tau)$ [12–14], is the intensity analogue of the first order function and measures the photon-number correlation. The function $g^{(2)}(\tau)$ can be written in terms of the number of counts on a photon detector, with the following expression

$$g^{(2)}(\tau) = \frac{\langle n_1(t)n_2(t + \tau) \rangle}{\langle n_1(t) \rangle \langle n_2(t + \tau) \rangle}.$$ (2.8)

Here, $n_{1,2}(t)$ are the number of photon clicks on detectors one and two at time, $t$. The function $g^{(2)}(\tau)$ can be interpreted as the probability of detecting a photon at time, $\tau$, after the previous photon detection. Figure 2.2a) schematically shows a typical experiment to measure $g^{(2)}(\tau)$. A HBT setup is used with SPDs at two outputs of a 50:50 BS. A delay, $\delta \tau$, is varied in one arm, while counting time resolved detections at SPD A and B. Light sources can be characterised into three categories

- bunched (super-Poisson statistics): $g^{(2)}(0) = 2$, $g^{(2)}(\tau) \to 1$ after a time longer than the coherence length, $\tau_c$.

- coherent (Poissonian/random): $g^{(2)}(0) = 1$ for a coherent source, and $g^{(2)}(0) \to 0$ for an antibunched source.

- antibunched (Sub-Poissonian): $g^{(2)}(\tau) \to 0$ for an antibunched source.
• coherent (Poisson statistics): \( g^{(2)}(0) = 1 \),
• and anti-bunched (sub-Poisson statistics): \( g^{(2)}(0) \to 0 \).

Single-photon detection alone is not sufficient experimental evidence for the existence of single photon quanta. It is enough to use a semi-classical treatment and describe this effect with a classical light wave and the spontaneous emission of a quantised electronic state in the photo detectors. However, a measurement of \( g^{(2)}(0) = 0 \) is a completely non-classical effect, and the first strong piece of evidence for quantised light. New Zealand Professor, Dan Walls, was a key contributor in the search for evidence of the quantum nature of light [2]. His seminal theoretical paper on anti-bunched light from the spectrum of resonance fluorescence was a key development in the study of non-classical light.

2.1.2 The single photon

The previous section demonstrated how the second-order correlation, \( g^{(2)}(\tau) \), can experimentally indicate the non-classical nature of light. If it is accepted that light is quantised, how can it be described? An ideal single photon should have a specific polarisation that can be linear, elliptical or circular, with the latter two embodying photon spin angular momentum. The spatial state can be defined by the waveguide mode or any superposition of modes, such as an orbital angular momentum mode. The temporal/spectral profile should be Fourier-limited. The formalism of first quantisation treats particles individually unnecessarily keeping track of each individual quantum particle. This is overly cumbersome; thus, this study moved to the second quantisation formalism [15]. In this formalism a single photon is described by the wave function, \( |\psi\rangle \),

\[
|\psi_{\text{photon}}\rangle = \int_0^\infty d\omega f(\omega)\hat{a}_\omega \hat{a}_\omega^\dagger |\text{vac}\rangle ,
\]  

(2.9)

where \( \hat{a}_\omega \) is the creation operator, acting on the vacuum, \( |\text{vac}\rangle \), to create a single photon, with a frequency distribution \( f(\omega) \) [16]. This formalism follows naturally from the concept of indistinguishability. If two photons have the same frequency, polarisation and mode then they are indistinguishable from each other. For indistinguishable photons, equation 2.9 reduces to \( \psi_n = |1\rangle \) a single photon in the occupation-number representation.
Moving into the photon population number picture, the number of photons in a state is represented by the photon-number state $|n\rangle$. The annihilation, $\hat{a}$, and creation, $\hat{a}^\dagger$, of a photon is described by the following operators

$$\hat{a} |n\rangle = \sqrt{n} |n-1\rangle \quad \text{and} \quad \hat{a}^\dagger |n\rangle = \sqrt{n+1} |n+1\rangle.$$ 

(2.10)

The photon number states can also be represented with creation operators acting on the vacuum state, $|\text{vac}\rangle$,

$$|0\rangle = \hat{a} |1\rangle = |\text{vac}\rangle \quad \text{and} \quad |1\rangle = \hat{a}^\dagger |\text{vac}\rangle. \quad \text{(2.11)}$$

**One-photon interference**

When a single photon is incident onto a 50:50 BS the wavefunction $|\Psi\rangle = |1\rangle_{a}$ is split into a superposition with half the probability amplitude being transmitted and the other reflected. Figure 2.3a) shows a Mach-Zehnder interferometers (MZI) with input and output port pairs labelled by the subscripts $a,b$ and $1,2$ respectively. The BS can be described by the unitary operator

$$\hat{U}_{BS} = \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 1 \\ 1 & -1 \end{bmatrix},$$

(2.12)

in the orthogonal basis set

$$|1\rangle_{a} |0\rangle_{b} = \begin{bmatrix} 1 \\ 0 \end{bmatrix}, \quad |0\rangle_{a} |1\rangle_{b} = \begin{bmatrix} 0 \\ 1 \end{bmatrix},$$

(2.13)

representing a photon entering either port $a$ or $b$, and the vacuum state at the other port. The output of the beam splitter is a linear superposition state, where

$$|1\rangle_{a} |0\rangle_{b} \xrightarrow{\hat{U}_{BS}} \frac{1}{\sqrt{2}} (|1\rangle_{c} |0\rangle_{d} + |0\rangle_{c} |1\rangle_{d}).$$

(2.14)

As the two parts of the superposition state pass through the upper and lower arms of the MZI, they acquire relative phase, described by the phase shifter
unitary operator

\[ \hat{U}_\phi = \begin{bmatrix} 1 & 0 \\ 0 & e^{-i\phi} \end{bmatrix}. \]  \hfill (2.15)

The second BS mixes the two states of the superposition, and the product of the phase shift and BS operation gives the MZI output state as

\[ \hat{U}_a \hat{U}_{\text{BS}} \rightarrow \frac{1}{2} (1 + e^{i\phi}) |1\rangle_A |0\rangle_B + \frac{1}{2} (1 - e^{i\phi}) |0\rangle_A |1\rangle_B. \]  \hfill (2.16)

This state yields oscillating probabilities

\[ P_A = \frac{1}{2} (1 + \cos \phi) \] 

and

\[ P_B = \frac{1}{2} (1 - \cos \phi) \]

that the photon will be detected at detectors \( D_A \) and \( D_B \), respectively. This result, derived using a quantum description, is the same as in the classical case. In the next section we will see how the interference of two-photons at a BS can only be described by quantum theory.

**Two-photon interference**

This section will introduce the Hong-Ou-Mandel (HOM) effect, first observed experimentally by Hong, Ou and Mandel in 1987 [17]. This effect is key to linear optical quantum computing and is discussed further in section 2.1.3. The experiment requires shining two photons on the inputs of a BS, and measuring photon coincidences across the outputs using SPDs, as illustrated in Fig. 2.3b).

![Figure 2.3](image_url)

**Figure 2.3:** a) The probability of detection a single photon input to a MZI at detectors \( A \) and \( B \) is proportional to the relative phase difference in the two arms \( \phi \), the same as for classical light. b) The experimental setup for a two-photon interference experiment. The delay \( \delta \tau \) tunes the relative arrival time at the BS. c) Coincident detections between SPDs \( A \) and \( B \) drop to zero when the two photons reach the BS simultaneously, with this shape called a HOM dip [17].
A photon incident on a 50:50 BS has an independent and equal probability of being transmitted or reflected. Classically, when two photons are sent to the input ports of a BS simultaneously, the presence of the second photon does not affect the other. In the quantum description, the input state $|1\rangle_a |1\rangle_b$ is transformed by the BS operator $\hat{U}_{BS}$ to give the output

$$
|1\rangle_a |1\rangle_b \xrightarrow{\hat{U}_{BS}} \frac{1}{2} (|1\rangle_c |0\rangle_d + |0\rangle_c |1\rangle_d) \otimes (|1\rangle_c |0\rangle_d - |0\rangle_c |1\rangle_d),
$$

$$
\rightarrow \frac{1}{2} (|1\rangle_c |0\rangle_d + |0\rangle_c |1\rangle_d - |0\rangle_c |1\rangle_d + |0\rangle_c |2\rangle_d),
$$

$$
\rightarrow \frac{1}{\sqrt{2}} (|2\rangle_c |0\rangle_d + |0\rangle_c |2\rangle_d).
$$

The phase shift on reflection from the BS is imparted onto the probability amplitude. The two indistinguishable cases, where both photons are transmitted or both reflected, have a relative phase of $-1$ and coherently cancel. The final state then only includes the terms when the two photons bunch, exiting randomly from either output port.

When the photons arrive at different times at the BS, the photons act independently, and the coincidence detections rate is high. As the relative delay is reduced, and the photons begin to overlap at the BS, the coincidence rate drops significantly ideally to zero. However, as sources of single-photons and detectors are imperfect, this experimentally only approaches zero. The characteristic dip in coincidence detections that occurs at zero relative delay, shown in Fig. 2.3c), is called a ‘HOM dip’. For a classical light source, the HOM dip visibility remains above 0.5. A HOM dip measurement is classified as non-classical if the reduction in coincidence rate (called the visibility) is greater than 50%.
2.1.3 Quantum information with single photons

Quantum information processing (QIP) is the art of computation, communication and storage of quantum states. Complex quantum processors that exceed the limit of a classical computer have yet to be realised. Photons interact weakly with matter, a property ideal for quantum communication. Thus, future QIP devices will almost certainly include photons as part of their architecture. The next section examines quantum computation and communication with single-photon qubits. Continuous variable schemes are omitted and a review of this topic can be found in reference [18]. This section also omits discussion of photonic quantum memories, where recent life-time measurements have exceeded tens of micro seconds for a few photons [19–21].

QIP technology

In 1982 Richard Feynman argued that a classical computer cannot efficiently simulate the physical world, which is inherently quantum mechanical, and proposed the idea of a ‘quantum computer’ [22]. The idea was to use a physical quantum system that is well understood and controllable to simulate and study an unknown quantum system. The typical choice for the quantum system is a particle with a binary pair of addressable states, called a ‘qubit’. A qubit-based quantum computer works by encoding information onto physical qubits, applying a series of unitary quantum evolutions and then make a measurement on specific qubits.

The proposal of key quantum algorithms such as a fast database search [23] and fast prime factoring [24] catalysed immense growth in the number of researchers interested in QIP. In 2000 David DiVincenzo formulated a set of criteria for assessing the potential of a physical system for use as a quantum computer [25]. The criteria requires:

1. A scalable physical system with well characterised qubits.
2. The ability to initialise the state of the qubits to a pure state.
3. Relevant decoherence times are longer than the gate operations.
5. Qubit specific measurement capabilities.
Currently there is a global race to find the physical system that best satisfies these criterion, and the potential to be scaled to the many thousands of qubits required for complex error-corrected computations. Some of the most successful platforms are summarised below.

<table>
<thead>
<tr>
<th>Physical system</th>
<th>Qubit degree of freedom</th>
</tr>
</thead>
<tbody>
<tr>
<td>Photon</td>
<td>Polarisation/number/time-bin</td>
</tr>
<tr>
<td>Electron</td>
<td>Spin/number</td>
</tr>
<tr>
<td>Superconductor junction</td>
<td>Charge/flux/phase</td>
</tr>
<tr>
<td>Trapped ion</td>
<td>Hyperfine energy levels</td>
</tr>
</tbody>
</table>

DiVincenzo also included an additional two criteria for quantum communication, requiring;

1. The ability to convert stationary to flying qubits.
2. The ability to faithfully transmit qubits between specific locations.

The last requirement represents the key condition for realising secure quantum communication. While each platform has unique advantages and technical challenges, no dominant technology has emerged.

**Photonic qubits**

A photonic qubit can be encoded, for example, in the polarisation degree of freedom with horizontal, $|H\rangle$, and vertical, $|V\rangle$, polarisation states [26]. An important property of qubits is they can be prepared in any superposition of their two orthogonal states, such as, in the $\frac{1}{\sqrt{2}}(|H\rangle + |V\rangle)$ state. A photonic qubit can also be encoding using two spatially separated beams, angular momentum, frequency and early-late time-bins. Each qubit representation is equivalent and can be represented on the Poincare sphere, with the polarisation example illustrated in Fig. 2.4. Rotation, $\theta$, in the vertical direction changes the relative superposition and rotations on the horizontal plane, $\phi$ the relative phase between $|H\rangle$ and $|V\rangle$, respectively.

The unitary evolution of a qubit state is described by a Hamilton function. The Hamiltonian forms an eigenvalue problem where the function applied to
\[ V + H = \frac{1}{\sqrt{2}} (|H\rangle - |V\rangle) \]
\[ V - H = -\frac{1}{\sqrt{2}} |\psi_{\text{photon}}\rangle \]

\[ |\rangle = \frac{1}{\sqrt{2}} (|H\rangle - |V\rangle) \]
\[ |+\rangle = \frac{1}{\sqrt{2}} (|H\rangle + |V\rangle) \]

**Figure 2.4: Qubit representation on the Poincare sphere.** The state of the photonic qubit \(|\psi_{\text{photon}}\rangle\) is represented by the red vector. When the vector points to the north or south pole, the polarisation state is \(|H\rangle\) or \(|V\rangle\), respectively. When the vector points elsewhere, it is in a complex superposition \(\alpha|H\rangle + \beta|V\rangle\).

the state gives the eigenvalues (energy) of each mode of the state. This is called the time independent Schrödinger equation and follows the form

\[ \hat{H} |\Psi\rangle = E |\Psi\rangle. \] (2.17)

An example of a Hamiltonian, key to optical quantum computing, is the beam splitter (BS). The BS Hamiltonian can be written as

\[ \hat{H}_{\text{BS}} = \theta \exp^{i\phi} \hat{a}_{\text{in}}^\dagger \hat{b}_{\text{in}} + \theta \exp^{-i\phi} \hat{a}_{\text{in}} \hat{b}_{\text{in}}^\dagger, \] (2.18)

where \(\hat{a}_{\text{in}}^\dagger, \hat{b}_{\text{in}}^\dagger\) are photon inputs for mode a and b of the BS, and \(\theta, \phi\) refer to the location of the qubit vector on the Poincare sphere [27].

**Nonlinear optical quantum computing**

The unitary evolutions required for a universal quantum computer are single-qubit and two-qubit operations. Single qubit gates act to rotate the qubit state. The ease of implementing single-qubit gates experimentally depends on the degree of freedom used for encoding. Two-qubit gates become a challenge because one photon is required to affect another. The property of weak interaction with matter, which makes photons desirable for quantum communication,
becomes a problem for multi-qubit computations.

It was first thought that strong optical nonlinearity, such as those found in a Kerr medium, was required to build an optical quantum computer [28]. For a successful two-photon gate, the control qubit must impart a large enough nonlinear phase shift onto a second target qubit. One possible nonlinear effect that was proposed for QIP is cross phase modulation (XPM), described by the Hamiltonian as

\[
H_{\text{Kerr}} = \kappa \hat{n}_C \hat{n}_T, \tag{2.19}
\]

where \( \kappa \) is the Kerr coefficient. \( \hat{n}_C \) and \( \hat{n}_T \) are the number operators for the control and target qubits, respectively. XPM induces a phase shift on the target qubit dependent on the number of control photons present. A controlled-phase (CZ) gate theoretically could be constructed using the cross-Kerr nonlinearity to couple the arms of two MZIs, with a magnitude that results in a \( \pi \) phase shift [29]. This proposal is shown schematically in Fig. 2.5a). A CZ-gate in conjunction with single-qubit gates forms a universal quantum computer.

It became clear in the late 1990s that achieving the required \( \pi \) nonlinear phase shift involving only two photons, while theoretically possible, was too challenging to implement experimentally [30–32]. Further, it was shown theoretically in 2006 that the non-instantaneous nature of the \( \chi^{(3)} \) nonlinearity would prevent high-fidelity quantum gates from ever being realised.

Despite the experimental challenges and theoretical forewarning, recent demonstrations have begun to rekindle the idea of quantum computing via photon-photon nonlinearity. Some key results include the use of classical light

\[50:50\]

\[a) \text{Cross-Kerr Nonlinear C-Phase gate}\]

\[b) \text{Linear optics Knill (C-Phase) gate}\]

**Figure 2.5:** a) A Kerr nonlinearity based C-Phase gate [29]. b) A linear C-Phase (Knill) gate, designed by E. Knill [33].
to impart a nonlinear phase shift to a single photon [34–36]. A photonic crystal photon number dependent switch has been demonstrated [37], and, recently, ultrahigh Q photonic crystal cavities were fabricated, opening the way for further development [38]. Experimental single-photon Kerr nonlinearities have been proposed and demonstrated using Rydberg atoms [39–42]. In addition to this, a scheme for optical QIP using Rydberg gates was recently proposed that potentially circumvents the limitations to the fidelity of nonlinear quantum gates [43]. This very recent work with Rydberg atoms is a particularly promising route towards realising a nonlinear optical quantum processor.

Linear optical quantum computing (LOQC)

In 2001, E. Knill, R. Laflamme and G. J. Milburn proposed an efficient photonic QIP scheme requiring only linear optics [44]. Using linear optics means that, ideally, the photon number, $|n\rangle$, is preserved. The scheme is efficient because the resources needed scale at a polynomial rate, in contrast to the previous LOQC schemes that required exponential resources [45]. The Knill-Laflamme-Milburne (KLM) proposal is based on single-photon sources, beam splitters, phase shifters, photon detectors and feedback from detection outputs.

The KLM proposal moves the nonlinearity from the photon interaction, to the measurement. The SPD click response to a photon detection is nonlinear [44]. This is not enough to create a deterministic quantum gate, but a probabilistic set of gates can be realised. The HOM interference of two single photons, central to the KLM scheme, has already been experimentally with high efficiency [5].

The linear version of the CZ gate, illustrated in Fig. 2.5b), contains only beamsplitters and phase shifters, and relies on the detection of ancillary photons to verify the success of the gate [33]. The most efficient LOQC CZ-gate has a probability of success of $2/27$. Feed-forward detection is an important part of the KLM LOQC scheme that can be used to increase the efficiency of gate operations [46]. As LOQC relies on a complex cascade of interferometers, interferometric stability becomes a key experimental factor. This is one of the key motivations of moving to photonic chip technologies of QIP, and is discussed in more detail in Section 2.3.
Quantum cryptography with single photons

Quantum cryptography (QC) has become an exciting real-world application with multi-disciplinary interest [47]. QC uses the indivisibility of a photon and entanglement for secure random key distribution. QC schemes are broken into three broad categories: those that require a single-photon, entangled photon-pairs or a coherent state (known as a continuous variable scheme).

In 1984, Bennett and Brassard proposed the first QC scheme which was based on encoding a single photon, known now the ‘BB84 scheme’ [48–50], illustrated in Fig. 2.6. Typically, polarisation is used as the encoding basis; however, the arrival time and orbital angular momentum can also be used. In the BB84 scheme, two different basis sets of pairs of orthogonal polarisations are chosen, such as $|H\rangle$, $|V\rangle$ and $|D\rangle$, $|A\rangle$ (diagonal/anti-diagonal). The sender A and receiver B randomly switch the basis in which they send and receive. After a photon is sent and received, A and B compare their basis choices. If they are the same, the measurement is kept. If an eavesdropper measures the photon in the wrong basis, the photon will be projected into an incorrect basis state, and B will detect an increase in the error rate.

Fibre optics provide the ideal channel for transmitting photons, via piggy-backing on the current global fibre optic network. Long-distance single-photon transmission is currently limited by losses to $\sim 200$ km [51–53]. As a quantum state cannot be copied; therefore, a classical amplifier cannot be used [47]. However, there are several proposal to extend the maximum distance including satellite links [51,54–56], quantum repeaters [57,58] and quantum relays [59].

![Figure 2.6: Schematic of the BB84 scheme. User A randomly encodes sent photons in one of two different basis sets. User B randomly measures in the same two basis. Both users then broadcast their basis choices after the photons are sent, keeping events only when the same choice is made. An eavesdropper will inevitably choose the wrong basis sometimes, thereby projecting the photon into the incorrect basis and, causing an increase in channel errors.](image)
2.2 Nonlinear optics for single-photon generation

The previous section, established the importance of developing high-quality sources of single photons. This section now turns to the methods by which this might be accomplished, and presents a realistic description of what is possible today.

2.2.1 The challenge of single-photon generation

The ideal single photon should typically have a specified spatial and polarization state, and a temporal/spectral profile that is near Fourier-transform limited. It is also desired that they are generated on-demand, and that each successive photon is indistinguishable to the previous. Currently, it is not possible to create a single photon with a high certainty without also having a high probability of generating two or more photons. This makes it very challenging to obtain a single photon in a truly on-demand manner. Generating single photons when and how they are required is a fundamental challenge in quantum optics. The non-deterministic nature of photon generation limits the feasibility of next-generation quantum processors that will require multiple simultaneous single-photon inputs.

There are three methods typically used to create single photons. First, by exciting an atom-like system, a single photon is emitted as it decays to the ground state. Examples of such systems include quantum dots [60,61] and nitrogen vacancy colour centres [62,63]. The challenge here is that the systems available often come with cryogenic requirements or poor emission directivity.

A single-photon source can be created by almost completely attenuating a laser, thereby leaving a small probability that a single photon will reach the

![Figure 2.7: Schematic illustration for the generation of photons using atom–like sources. Imperfect collection leads to a stochastic photon output.](image-url)
output in a given time. The likelihood of two or more photons reaching the output is even lower. The ratio of single to multiple-photon probabilities is fixed by Poisson statistics, thereby limiting the usefulness of these sources.

The last method is photon pair generation via a parametric nonlinear process, such as spontaneous four-wave mixing (SFWM) or down-conversion. The challenge here again lies in the intrinsic link between the rate of useful single-photon creation, and how often two or more photons are generated. The ratio of single-pair to multi-pair photon probabilities is fixed by Poisson or thermal statistics \[64\], with the probability of \(i\) photon-pairs given by

\[
P_{\text{Poisson}}(\mu, i) \sim e^{-\mu} \frac{\mu^i}{i!},
\]

\[
P_{\text{Thermal}}(\mu, i) \sim \frac{\mu^i}{(1 + \mu)^{i+1}}.
\]

Here, \(\mu\) is the average number of pairs generated per pump pulse. By expanding out, for example, the Poisson series the probability that a given photon pair number is generated can be calculated \(\sim e^{-\mu} |0\rangle_{\text{pair}} + e^{-\mu} \mu |1\rangle_{\text{pair}} + e^{-\mu} \frac{\mu^2}{2} |2\rangle_{\text{pair}} + ...\) and so on. Whether the photon pair number distribution follows a thermal or Poissonian distribution depends on the ratio of the length of the pump pulse \(T_{\text{pump}}\) over the coherence length of the photon pairs \(\tau_{\text{photon}}\). For a Poisson source, \(T_{pump}/\tau_{\text{photon}} \gg 1\), with photon pairs generated in a number of temporal modes. In contrast a thermal source has \(T_{\text{pump}}/\tau_{\text{photon}} = 1\) and all the photons are generated coherently in a single temporal mode. Details of when the photon number distribution follows thermal or Poissonian statistics are presented in reference \[64\].

In all three cases, generating higher rates of single photons is accompanied
by a higher proportion of unwanted additional photons. However, photon-pairs generated by nonlinear processes do have an advantage. Detecting one photon from a pair indicates the existence and timing of the other photon—known as ‘heralded single-photon generation’. This added knowledge makes it possible to deterministically combine the output of several spatially separated sources, or combine multiple temporal modes of a single source. This results in a higher single-photon rate without the corresponding increase in multi-photons, thereby breaking the intrinsic link between the photon-pair number probabilities. These multiplexing schemes for improving photon source statistics are promising, and are explained in detail in Section 2.2.5.

### 2.2.2 Photon-pair generation in nonlinear media

When an EM field interacts with a dielectric medium, the internal electronic distribution of the dielectric is polarised. For low field intensities, the electronic response is linear, that is no new frequencies appear. For higher field intensities, the electronic response behaves nonlinearly, and the induced polarisation, $\tilde{P}$, is expressed by

$$\tilde{P} = \epsilon_0 (\chi^{(1)} E_1 + \chi^{(2)} E_1 E_2 + \chi^{(3)} E_1 E_2 E_3 + \ldots),$$

(2.22)

where $E_i$ is the electric field, $\epsilon_0$ is the vacuum permittivity, $\chi^{(1)}$ is the linear susceptibility, and $\chi^{(2)}$, $\chi^{(3)}$ are the second and third order nonlinear susceptibilities, respectively [65,66]. EM fields of different frequencies can interact via $\chi^{(2)}$, $\chi^{(3)}$ nonlinearities, thereby allowing frequency conversion. The allowed frequencies are restricted, initially by energy and momentum conservation.
The $\chi^{(2)}$ nonlinearity has primarily two broad categories of effects. The first two fields combine to generate a single new field, which includes the classical effects of sum and difference frequency generation. In the second effect, a single EM field generates two new fields at different frequencies, called ‘spontaneous parametric down conversion’ (SPDC). This is a spontaneous quantum process in which a pump photon $\omega_p$ is probabilistically converted into a correlated photon-pair $\omega_s$ and $\omega_i$ where $\omega_p = \omega_s + \omega_i$ [67], as shown in Fig. 2.10. The $\chi^{(2)}$ nonlinearity is only naturally present in material with a non-centro-symmetric crystal lattice.

The $\chi^{(3)}$ nonlinearity has a variety of associated classical nonlinear effects. These can be broken into roughly two groups—those associated with the instantaneous Kerr nonlinearity, and those with a slower response. While the $\chi^{(3)}$ effects are weaker than the $\chi^{(2)}$ effects, the $\chi^{(3)}$ nonlinearity is present in all materials, including both amorphous and crystalline materials, and both glass and semiconductors. Classical $\chi^{(3)}$ frequency conversion is most dramatically observed via stimulated four-wave mixing. Here, the wavelength of a seed field is converted to another wavelength by either a pair of degenerate or non-degenerate pumps. SFWM is an essentially identical process, with the role of the seed field taken by random quantum fluctuations. Like SPDC, SFWM is a spontaneous process, where two pump photons, $\omega_p$, are converted into a correlated photon-pair with energy conservation,

$$2\omega_p = \omega_s + \omega_i,$$  \hspace{1cm} (2.23)
and phase matching condition,

$$\kappa = k_s + k_i - k_{p1} - k_{p2} + \gamma (P_{p1} + P_{p2}),$$ (2.24)

where $\kappa$ is momentum phase mismatch; $k$ is the wavevector of the photons; $P_{p1} + P_{p2}$ are the power of the pump fields (reducing to $P_{\text{pump}}$ in the degenerate case of $P_{p1} = P_{p2}$); $\gamma$ is the medium nonlinear coefficient; and $\gamma = n_2 \omega / (c A_{\text{eff}})$, where $n_2$ is the nonlinear refractive index at frequency, $\omega$, and $A_{\text{eff}}$ is the effective mode area of pump field when it is confined to a waveguide. The $\gamma (P_{p1} + P_{p2})$ term appears as a result of the induced phase mismatch resulting from the Kerr nonlinear effects of self and cross phase modulation [66,68], both result in an intensity dependant phase shift.

**Photon statistics**

Photon-pair sources have yet to be ideally implemented. Many practical considerations degrade the performance of a pair source, such as coupling and propagation loss, and detector and collection efficiencies. Ideally, to completely characterise a photon source, it is necessary to know the exact photon statistics of the photon pairs generated. To do this, one needs unit efficiency photon-number resolving detectors. This gives a direct measurement of the source’s photon statistics; however, to achieve this experimentally is very challenging. Alternatively it is possible to use inefficient and non-number resolving SPDs to indirectly investigate the source characteristics by measuring the the coincidence to accidental ratio (CAR).

Figure 2.11 lays out a typical photon pair generation measurement. The convention ‘signal’ and ‘idler’ is typically used to label the higher and lower energy photon from a pair, respectively. A laser pumps a nonlinear waveguide to generate pairs of photons, and the pump is dropped and idler and signal photons wavelength separated, before being detected, and the photon detections are time tagged by the electronic time interval analyser. Correlated detections on the two SPDs are called ‘coincidences’ $C$ (illustrated by the orange dashed oval in Fig. 2.11). The portion of these that are actually photon pairs generated in the nonlinear waveguide are called ‘true coincidences’ $C_{\text{true}}$. 

24
Figure 2.11: Schematic of a typical photon pair source characterisation setup. A laser pulse generates photon pairs in a nonlinear waveguide. The pump is dropped and the photons’ wavelength is divided before reaching a pair of SPDs. Photon arrival times are used to calculate the photon statistics.

Accidentals, $A$, are detections correlated across one or more pump pulse periods, $T$, as illustrated by the green dashed oval in Fig. 2.11. These are due to any combination of multi-pair generations, SpRS photons, detector dark counts or leaked pump photons. The CAR is the ratio of the true coincidences to accidental coincidences

$$\text{CAR} = \frac{C_{\text{true}}}{A} = \frac{C - A}{A}. \quad (2.25)$$

The CAR can be derived theoretically to different levels of completeness, each giving insight to the physical processes involved [64]. Here $\mu$ again represents the average number of photon pairs generated per pump pulse. If we first neglect uncorrelated events, and assume $\mu << 1$, the probability of a photon-pair detection per pump pulse is $\mu \eta_i \eta_s$, where $\eta_i/s$ is the collection efficiency in the idler and signal channels, respectively. Multiplying this probability by the laser pulse repetition rate, $R$, gives the true coincidence-rate $C_{\text{true}} \sim \mu \eta_i \eta_s R$. Similarly the probability of a single photon detection at each channel is $\mu \eta_i/s$. Thus, the accidentals rate $A$ is the product of the probability of a photon detection at each channel and the laser repetition rate, giving $A \sim \mu \eta_i \mu \eta_s R$. In this noiseless approximation, the $\text{CAR} \sim 1/\mu$, as shown by the dashed line in Fig. 2.12. The CAR is reduced by multi-pair generation for increasing $\mu$.

Now, when uncorrelated events, such as dark counts, are included the CAR curves appear qualitatively much closer to the experimental data. The single counts are now $N_x = \mu \eta_x + d_x$, where $d_x$ is the dark counts in channel $x$. This
Figure 2.12: Theoretical CAR curves. The blue dashed curve follows a $1/\mu$ trend, with no dark counts. The yellow curve has dark counts, and the red curve has equal dark counts with twice the loss.

gives an expression for the CAR, where

$$CAR = \frac{\mu \eta_s \eta_i}{(\mu \eta_s + d_s)(\mu \eta_i + d_i)}. \quad (2.26)$$

As $\mu \to 0$, the CAR quickly drops to zero, as the dark counts begin to dominate over real photon detection counts. Equation 2.26 already includes enough terms to provide qualitative intuition of experimental data. For a more complete CAR expression that takes into account higher orders of multi-pair generation and additional linear noise sources ($\propto \sqrt{\mu}$), such as Raman and pump leakage, see [69,70] and Section 4.3.2.

Nonlinear photon-pair generation

Building a high-quality integrated source of single photons remains a key challenge. Various platforms are under investigation, including on-demand sources such as quantum dots [71] and diamond vacancy centres [72] as well as probabilistic photon-pair generation in nonlinear waveguides. On-demand sources often require complex experimental techniques, such as liquid helium cooling and confocal microscopy, and producing samples that emit in the desired near-infrared remains challenging [73,74]. Nonlinear waveguides have emerged, thus far, as the choice for quantum computation and communication research due to their room temperature operation and directed photon output. Photon pairs were first generated in bulk nonlinear crystals in the 1980s for photon-based
tests of Bell inequalities [75]. Improved versions of these bulk sources were used in the 1990s for the first demonstrations of quantum state teleportation [76] and entangled photon generation [77]. The 1990s also saw rapid growth in the field of guided wave nonlinear optics. This created the tools required to use photonic waveguides for photon pair generation.

The first waveguide photon generation experiment was performed in 1999 using a periodically poled silica fibre to enhance the $\chi^{(2)}$ nonlinearity for efficient SPDC [78]. In 2001 the first on-chip SPDC experiment was performed in periodically poled lithium niobate (PPLN) [79]. This was followed that same year by the first $\chi^{(3)}$ fibre source [80], which had a narrow SFWM output spectrum. In 2005 a dispersion engineered photonic crystal fibre was used to create photon pairs, by SFWM, with a larger frequency separation [81].

To achieve further integration a platform with a high refractive index and mature fabrication technology was needed. Piggy-backing on the semiconductor industry, planar silicon waveguide circuits have been fabricated with unprecedented device density. In 2006 photon pair generation in silicon was both proposed [82] and demonstrated [83]. The next year a silicon entangled pair source was demonstrated [84] and in 2011 photon pairs were generated in an ultra-compact 100 $\mu$m-long slow-light photonic crystal device. With characteristic sizes on the order of 100 $\mu$m and silicon wafer diameters on the order of 100 mm, scaling to complex quantum photonic circuits is now possible.

However, chip-scale photon sources, whether they are based on $\chi^{(2)}$ materials (such as PPLN) or $\chi^{(3)}$ materials (such as silicon), have some inherent drawbacks. PPLN waveguides require bulky temperature control, while silicon suffers from TPA and associated free-carrier effects that together cause high nonlinear loss. There is also significant interest in using amorphous glass waveguides, such as chalcogenide for chip-scale photon-pair generation [85,86]; entanglement generation [87,88]; and the frequency conversion of single photons [89,90]. However, SpRS is a detrimental noise source in such glass devices [86]. Partial suppression of SpRS has been demonstrated for silica glass using cryogenic cooling [91–95].

The next sections review the current state of the art in integrated nonlinear photon generation in amorphous glass and crystalline silicon.
2.2.3 Nonlinear photon generation in glass

Nonlinear optics in waveguides were born out of two key technological inventions. The laser and the optical fibre. Coherent and high power laser light in combination with the small core area, long length and low loss of modern optical fibre enabled high nonlinearities to be achieved.

The first experiments aiming to generate photons using nonlinearity naturally considered amorphous glass fibres as a starting point. Eventually, further integration was desired to gain improved stability and scalability. This led to the exploration of silica direct write waveguides and highly nonlinear chalcogenide waveguides. The challenges that remain for amorphous glass quantum photonic devices are: (i) uncorrelated noise due to Raman scattering and (ii) further characteristic size reduction.

Spontaneous Raman scattering

Raman scattering occurs when photons interact with the phonon bath in a material, where phonons are the quanta of molecular vibration. A photon interacting with any media has a finite probability of losing energy to a phonon mode in the bath (Stokes scattering, see Fig. 2.13a), or absorbing energy from a phonon mode (anti-Stokes scattering, Fig. 2.13b). Stokes and anti-Stokes scattering correspond to a red and blue shift of the photon in wavelength, respectively.

Figure 2.13: When photons from the pump laser interact with the molecule of a material, they can transfer/absorb energy to a phonon (vibrational) mode, corresponding to a) Stokes scattering and b) anti-Stokes scattering.
In the weak pump power regime Raman scattering is spontaneous (SpRS) and proportional to the density of phonon states in the material and to the thermal occupation of the phonon bath. SpRS scales linearly with the length of the sample and the pump power. In the high pump power regime Stimulated Raman scattering can occur in long waveguides leading to exponential nonlinear gain. For more details on the theoretical description of SpRS, refer to section 3.2.2. In amorphous material, SpRS is broadband and can overlap with the SFWM spectrum. This introduces uncorrelated photon noise that reduces the CAR performance of amorphous photon-pair sources.

**SFW in optical fibres**

The natural first choice for photon pair generation in amorphous media was in silica glass fibres. These were already becoming the backbone of the telecommunication industry, and the fabrication of low loss fibres rapidly advanced. Initial experiments were limited to anomalous dispersion fibres, with SFWM photons generated close about the pump frequency [80, 96–98]. These initial demonstrations were limited by uncorrelated SpRS noise generated in the SFWM bands. This was partly countered by cryogenic cooling in liquid nitrogen and helium [92–94], as well as operating at very small detunings [91,98].

Photonic crystal fibres (PCFs) can be either index guiding or guide using a Bragg mirror realised with a periodic array of air holes in the fibre. PCFs allow tailored dispersion and birefringence, thereby enabling SFWM pair photons to be generated at large detuning (typically > 70 THz) from the pump, beyond the SpRS bandwidth, which results in improved CAR performance [81,99–104].

![Amorphous glass devices](image)

Figure 2.14: Amorphous glass devices: a) Optical fibre. b) Photonic crystal fibre [104]. c) Silica direct write waveguide [105]. d) Chalcogenide chip.
Silica direct write waveguides (DRWs)

Waveguide circuits can be created in solid silica chips for both pair generation and LOQC gates. An example of a $3 \times 3$ multiport interferometer, sometimes known as a tritter, is shown in Fig. 2.14c. The fabrication of silica waveguides is undertaken using direct laser writing. A femtosecond pulse in the visible wavelength range is focused into the glass with enough energy to modify the refractive index and create a waveguide. These pulses are typically generated using a titanium-sapphire laser at a repetition rate ranging from kilohertz to megahertz. The waveguide profile is innately elliptical due to the elongation of the focal spot inside the glass. This can be corrected for, however elliptical waveguides are birefringent, and this can be a desired property for SFWM where the photon pairs are far detuned in wavelength. For example, by pumping a 4 cm-long birefringent silica waveguide at 729 nm signal and idler photons were able to be generated efficiently at 676 nm and 790 nm, respectively [106]. This is relatively long for an integrated device due to the low intrinsic nonlinearity of silica. The advantage of DRWs is that they are well matched to silica fibre for efficient out-coupling of the generated photon pairs.

Chalcogenide waveguides

Chalcogenide [107] has been a flagship photonic material for the last two decades, especially in the Australian context, where it has historically been the key platform in the Centre for Ultrahigh bandwidth Devices for Optical Systems (CUDOS). Chalcogenide is transparent from $\sim 800$ nm well into the mid infrared (MIR). The nonlinearity is $100 - 1000$ times greater than silica, and exhibits little nonlinear loss in the form of two photon absorption. The waveguide geometry can be tuned to create dispersion engineered waveguides and, as it is photo refractive, it can be post-processed using laser writing to create Bragg gratings [108]. Chalcogenide has been successfully applied to nonlinear optics [109], chemical sensing [110, 111], distributed optical fibre temperature sensors [112], on-chip radio frequency spectrum analysis [113] and MIR spanning supercontinuum generation [114]. The maturity and diversity of the platform make it very attractive for quantum applications.
In early 2011, the team at CUDOS investigated photon-pair generation at 1,550 nm using SFWM [86] in As$_2$S$_3$ chalcogenide. A correlation peak was observed; however, it was nested in a sea of uncorrelated noise photon counts. The raw data from this measurement are replotted in Fig. 2.15a). The CAR (see Section 2.2.2) was measured using a coincidence detection scheme, with the results shown in Fig. 2.15b). The authors claim that the contribution from SpRS dominated the contribution of the SFWM-correlated pair generation, despite the SFWM efficiency being high in the detuning regime explored [85].

In addition, out-coupling photon pairs from a chalcogenide waveguide ($n = 2.4$) to a silica fibre ($n = 1.5$) remains a significant point of loss.

### 2.2.4 Nonlinear photon generation in silicon

Silicon has a very high refractive index ($n \sim 3.4$), allowing single mode waveguides to be made with a cross-sectional area of ~0.1 $\mu m^2$ and waveguide bend radius of a few micrometres. Due to its compatibility with the electronics industry, complementary metal-oxide-semiconductor (CMOS) standard fabrication provides a mature toolset for creating highly complex and mass producible devices. Typically high quality, epitaxially grown crystalline silicon-on-insulator (SOI) wafers are used for photonic devices. While the discussion following in this chapter will be limited to crystalline silicon, there have been
demonstrations of nonlinear photon generation in amorphous silicon, which suffers from Raman noise similarly to chalcogenide [115].

The challenges remaining for silicon quantum photonics are the relatively high propagation and coupling losses, as well as nonlinear loss from two-photon absorption. The propagation loss is due to the larger role played by surface roughness for smaller waveguides. Similarly, the mode area mismatch with silica fibre is large, which is a challenge when coupling light in and out of silicon waveguides. This section reviews the progress in silicon photon generation.

Two photon absorption (TPA)

By definition, semiconductor material has an electronic band gap between the insulating and conducting bands. For a band gap with energy $2X_eV$, two photons of energy $X$ can be absorbed simultaneously in order, to eject an electron from the insulating band and create a free electron, or free carrier. These free

![Diagram of Silicon and GaInP electronic band diagrams with TPA and TPA-free coincidence experiments.](image)

Figure 2.16: a) Silicon electronic band diagram, with photon pair coincidence experiment showing limiting effect of TPA. b) GaInP electronic band diagram, with TPA-free photon pair coincidence experiment [70].
carriers both absorb (free carrier absorption – FCA) light and modify the refractive index of the material (free carrier dispersion – FCD). The timescale for free-carrier dynamics is typically around a nanosecond, with measured carrier lifetimes on the order of 0.5 ns [116]. For weak pulse powers used in nonlinear photon generation, the effects of FCA and FCD can be neglected to focus on TPA [117].

The main effect of TPA is to cap the peak power in the waveguide [116], thereby limiting the nonlinear photon generation rate [70]. Silicon has a small (1.2eV) indirect band gap. Here, indirect means that to satisfy momentum conservation, the TPA must be phonon assisted, as illustrated in Fig. 2.16a). Photons with a wavelength below $2\mu m$ have enough energy for TPA. This limits the peak power of photon pairs generated around 1550nm, as shown in Fig. 2.16a). III/V semiconductors, such as GaInP, have a larger band gap (1.8eV) and photons with a wavelength of above 1380nm (including the telecom bands) do not have enough energy for efficient TPA. For GaInP, it is clear that nonlinear photon pair generation is uninhibited by TPA absorption [70, 118]. It is worth noting that three-photon absorption is observed in GaInP if the input power is increased further, and that this is generally above the range of interest for photon pair generation.

Silicon nanowires

In 2006, there was a paradigm shift with the first demonstrations of SFWM correlated photon pair generation in a silicon nanowire [83]. Now photon pairs could be generated on the millimetre scale, integrated with complex LOQC circuits, and fabricated on scale using mature CMOS fabrication technology.

![Figure 2.17: Silicon devices: a) Nanowire [119], b) Ring resonator [120], c) Microdisk [121] and d) Photonic crystal waveguide [86].](image)
Silicon was also the first platform where SFWM, with a small pump detuning, was shown with low noise [122], owing to the narrowband Raman spectrum in silicon. Initially, the high coupling loss from a silicon waveguide to an optical fibre limited the brightness of these sources. Inverse tapers and polymer mode converters have been implemented to reduce these losses, thereby increasing the number of photons that make it off the chip. In 2011 the Nippon Telegraph and Telephone (NTT) basic research laboratory was able to show a HOM interference dip between two independent bright silicon photon pair sources, thereby indicating their indistinguishability [123].

**Silicon resonant structures**

In a micro-ring or micro-disk structure, photons can travel along a closed loop for many orbits; thus, the nonlinear interaction between photons and the device can be enhanced. The nanowire dimension determines the basic SFWM properties such as phase-matching bandwidth, while the ring circumference together with the ring-bus gap determines the SFWM enhancement.

Silicon microdisks sit on a silica pedestal, with air above and below for most of the disk. The sharp contrast in refractive index of air and silicon strongly confines the optical modes. Photon pairs are coupled evanescently to and from the disk with a tapered silica fibre [121]. The resonant enhancement to the SFWM and efficient coupling has enabled CAR values as high as 1000 in these structures. Further, this has been demonstrated at pump powers of up to five orders of magnitude lower than that required for a standard silicon nanowire [120,121].

A coupled resonator optical waveguide (CROW) consists of a one dimensional (1D) array of coupled resonators [124], that provide slow-light enhancements [125]. They have a wider pass-band than a single resonator structure, and have been implemented using rings for efficient photon pair generation [126,127]. Photonic crystal cavities can also be used as the base element of a CROW, and were recently employed in a demonstration of nonlinear photon pair generation in silicon [128].
Slow-light silicon photonic crystal waveguides

To achieve truly scalable photonic quantum devices, the footprint of nonlinear photon pair sources needs to be minimal. As in the previous section, with the exception of resonant structures, the characteristic length of a device is one centimetre (silicon nanowire) or several centimetres (chalcogenide, silica). Photonic crystal waveguides (PhCWGs) have been shown to drastically enhance the nonlinearity of a silicon waveguide through slow-light propagation (see Section 2.3.3). Dispersion engineering has allowed broadband stimulated four-wave mixing in compact 100 µm devices [129, 130].

In 2011, the team at CUDOS demonstrated correlated nonlinear photon-pair generation via SFWM in a 96 µm long silicon PhCWG [86]. The device was dispersion engineered to have a slow-light band over a bandwidth of 15 nm, centred at 1,553 nm, close to the middle of the telecommunication C-band. The most relevant property that is ‘slowed’ is the group velocity $v_g$, defined as the speed at which a pulse envelope propagates [125], which in this band is $\sim \frac{c}{30}$. The dispersion and transmission profile of the waveguide is shown in Fig. 2.18a). The pump (green) is placed in the centre of the flat band, and the photon pairs collected at 4.8 nm detuning either side of the pump. The nonlinearity parameter $\gamma \sim 4,000/W/m$. At an input pulse, peak power of 0.24 W and generation probability of 0.006 pairs per pulse, the maximum measured CAR was 12.8. This is above the useful threshold of a CAR of 10, such as for use for secure quantum communication [92]. Figure 2.18b) shows the CAR and coincidence counts for the PhCWG. The coincidences

![Figure 2.18: Slow-light silicon PhCWG photon pair generation [86].](image)
initially follow a quadratic trend, as expected for SFWM, before rolling off as the effects of TPA limit the pump peak power. In subsequent work, the CAR was improved to 33 in 2012 [131] and 130 in 2013 [70]. Also in 2013, a TPA-free GaInP PhCWG was demonstrated with a CAR of 60, limited by linear loss [118].

2.2.5 Multiplexing: A route to on-demand photons

In Section (2.2.1), it was demonstrated that truly on-demand photon generation remains a challenge. As aforementioned, photon pair sources have the advantage over an attenuated laser that, by detecting one of the pair the existence and timing of the remaining photon is heralded. This is called a ‘heralded single-photon source’. An individual heralded single-photon source has fixed photon statistics that are either Poisson or thermal. However, by multiplexing multiple sources together, the intrinsic limits on the photon statistics can be broken. This section explains these approaches and discusses the demonstrations of their use to date.

Spatial multiplexing

Practically, this is undertaken by using the herald detections to reconfigure an $N \times 1$ low-loss optical switch, routing single-photons from multiple sources to a common single-mode output. This is shown in Fig. 2.19, where eight nonlinear photon pair sources (in this case, silicon PhC waveguides) have their outputs combined by an optical switch, triggered by a herald detection, to produce an enhanced single-photon output. The assumption is taken that we are operating in the limit of low multi-pair generation and low simultaneously pair generation in multiple sources. The probability of a single photon reaching the output of the multiplexing system is then

$$P_{\text{Spatial}}(\mu, N) \sim N \eta_S^{\log_2(N)} P_{\text{Source}}(\mu),$$

(2.27)

where $N$ is the number of multiplexed source ($N$ is a power of two), $\eta_s$ is the switch loss and $P_{\text{Source}}(\mu)$ is the probability for a single source. A more detailed expression for the multiplexing photon pair distribution can be found
in reference [16]. For large $N$ and zero loss ($\eta_s = 1$) and assuming perfect detector efficiency, the probability tends towards an on-demand single-photon source. In reality, there will always be loss and imperfect detection limiting the performance of any scheme. In addition, for a larger number of sources, the probability of two or more sources generating a single pair increases; thus, there are diminishing returns beyond a certain number of sources. The number of switches required to implement this scheme scales linearly, with $(N - 1)$ switches needed to multiplex $N$ sources.

The idea of spatial multiplexing was proposed by Alan Migdall et al. at NIST in 2002 [132]. In 2007, an alternate method was proposed that uses electro-optic polarisation controllers, where polarisation beam splitters are often lower loss than $2 \times 2$ Mach–Zehnder interferometric switches [133]. This scheme was first demonstrated by Xiao-song Ma et al., from the Zeilinger group at the University of Vienna, in 2011 [134]. Using a double pumped barium borate crystal and free-space polarisation switches, Ma et al. showed the potential for a four-fold increase in single-photon rate; however, the instability of bulk optics impeded scaling to more devices. Several additional publications have examined the limits of multiplexing [135, 136], and the practical requirements of using such a source for LOQC [44, 137].

The first demonstration of integrated spatial multiplexing was performed in late 2013 at the University of Sydney, with help from the Defence Science and Technology Organisation (DSTO) and Macquarie University. In this, two photonic crystal waveguide photon sources fabricated at the University of St Andrews and the University of York in the United Kingdom were combined on a single silicon chip [138, 139] using the Migdall scheme. Also in 2013, another alternate scheme with an asymmetric architecture for multiplexing was proposed. The researchers, theoretically propose that in certain regimes this scheme has potential to outperform the traditional symmetric switching architecture [140].

Very recently, in early 2014, a demonstration at Macquarie University, in collaboration with the University of Sydney and University of Nice Sophia Antipolis in France multiplexed four monolithic periodically poled lithium niobate (PPLN) waveguides, butt-coupled to direct laser-written silica wavelength di-
Temporal multiplexing

In the previous section, multiple sources spatially separated were combined using multiplexing. Each source has a single spatial mode, and these modes were combined deterministically using an $N \times 1$ optical switch. This section discusses temporal multiplexing where one or more photon pair source has many temporal modes actively combined by routing the photons through a different number of delay lines. Depending on the temporal mode in which a photon pair was created, as indicated by the timing of the herald detection, a different amount of relative delay is applied. This idea was proposed theoretically by J. Mower and D. Englund in 2011 [142], who envisaged the scheme being implemented with a sophisticated photonic integrated circuit.

In Fig. 2.20, a pulse train with period, $T$, pumps photon-pair generation
in a nonlinear device. The output is then divided, for this example, into four
time-bins, $t_i$, $i \in \{1, 4\}$. A herald detection, in time-bin $t_2$, for example,
will trigger the remaining single photon to be switched through $3T$ of relative
optical delay line moving it into time-bin $t_1$. For a herald detection in time-
bin $t_1$, zero relative delay is applied. The output has a period now of $4T$
with four times the probability of a single photon per clock cycle. A possible
setup for temporal multiplexing is shown in Fig. 2.21 with Mach-Zehnder $2 \times 2$
switches used to route the heralded photon though the $3T$ of relative delay. The
probability of a single photon reaching the output of a temporal multiplexing
system with $N$, ($N$ is a power of two), temporal modes is

$$P_{\text{Temporal}}(\mu, N) \sim N \eta N^{(\log_2 N + 1)} P_{\text{Source}}(\mu), \quad (2.28)$$

The number of switches required to implement this scheme scales logarith-
mically with $(\log_2 N + 1)$ switches needed to multiplex $N$ temporal modes of
the source. Thus, temporal multiplexing is more resource efficient than spa-
tial schemes for a high number of multiplexed temporal modes. However, it is
more challenging to implement because switching triggered on temporal modes

\[ t_1 \quad t_2 \quad t_3 \quad t_4 \]

\[ T \quad T \quad 2T \]

\[ \text{Source} \quad \text{FPGA} \quad \text{Dump} \]

\[ \text{MZ}_1 \quad \text{MZ}_2 \quad \text{MZ}_3 \]

\[ \text{Single photon detector} \]

\[ \text{Delayed single photon} \]

\[ \text{AWG} \]

**Figure 2.20:** Timing diagram of a four fold temporal multiplexing scheme.

**Figure 2.21:** Schematic illustration of temporal multiplexing of four temporal
modes using Mach-Zehnder optical switches (MZI). AWG = arrayed waveguide
grating.
is difficult to implement, and to date no experimental demonstration has been published.

In the future, this idea will need to be extended to combine many more devices onto a single optical chip. Undoubtedly, hybrid integration of several material platforms will be required to realise all of the components necessary, including delay lines and fast switches. Multiplexing will allow many single photons to be generated nearly on-demand, thereby enabling a new regime of complex quantum processing.
2.3 Photonic integration

As discussed in Section 2.1, quantum photonics is a modern field that emerged during the last 15 years. The vision forward is to use the mature tool box of photonics to create stable and scalable quantum information processors [4].

Quantum optics experiments are stochastic and must be repeated many times. The interference of two single photons on a beam splitter is a key part of linear optical quantum computers, (see Section 2.1.1). MZIs are often used to implement the beam splitter operation, especially in a planar geometry, making interferometric stability a vital property for scalable quantum circuits. The power of photonic integration can be illustrated using a simple calculation comparing the stability of a MZI made from silica fibre and an integrated silicon waveguide. A schematic of an MZI can be seen in Fig. 2.23.

Fringe visibility in a MZI is \( I_{vis} \propto (1+\cos\phi) \), where \( \phi = (2\pi L/\lambda_0)(dn/dT)\Delta T \) [143]. The temperature coefficient of refractive index at room temperature and with \( \lambda_0 = 1.5 \mu m \) light is \( 8 \times 10^{-6} T^{-1} \) for silica [144] and \( 2 \times 10^{-4} T^{-1} \) for silicon [145], respectively. Despite silicon having a higher coefficient, the characteristic length of a fused silica optical fibre MZI (~0.1 m) is much larger than an equivalent silicon waveguide MZI (~100 \( \mu \)m). For a 0.1 K temperature difference between the interferometer arms, the corresponding phase shifts are \( \Delta\phi_{SiO_2} \sim 0.1\pi \) for silica and \( \Delta\phi_{Si} \sim 0.003\pi \) for silicon. Here, a \( \pi \) phase shift corresponds to a change from a 50:50 splitter to complete constructive and de-

![Bulk optics](image1) ![Integrated optics](image2)

**Figure 2.22:** The quantum optics experiments traditionally realised in bulk optics suffer from challenges in stability and scalability. Moving to integrated photonics platforms moves quantum optics experiments into a new regime, leveraging the maturity of photonic fabrication.
Figure 2.23: a) Schematic of an MZI. b) Output intensity from Port 1 and Port 2 as a function of the phase different between the two MZI arms.

Structive interference at port 1 and 2, respectively. This means the silicon MZI is almost two orders of magnitude less sensitive to temperature fluctuations than the silica. This increase in stability results directly from the reduction in MZI size.

Further to this, smaller structures are likely to be closer to thermal equilibrium. For example, with a realistic laboratory temperature gradient of $\nabla T = 0.1\, K\, m^{-1}$, the temperature difference between the arms is $10^{-5}\, K$ for the silicon MZI and $0.1\, K$ for the silica fibre MZI. The phase shift is then $\Delta \phi_{Si} \sim 3 \times 10^{-7}\pi$, 40,000 times more stable in the silicon MZI.

### 2.3.1 Waveguides

Complex photonic circuits have been realised by chip-scale integration. Waveguides confine light to a localised spatial mode and enable precise control over the light path and phase. In quantum optics, control over the spatial mode is

Figure 2.24: a) Scanning electron micrograph of a silicon nanowire waveguide. The nanowire is 220nm high and 450nm wide silicon on a silica substrate [146]. b) Fundamental TE mode profile, simulated using an eigensolver [147].
important. A waveguide, when single moded, can provide automatic spatial overlap for two photons of the same frequency.

Light can be confined to a waveguide using several types of interfaces including index guiding (total internal reflection), metal mirrors and Bragg (dielectric) mirrors. A silicon nanowire is a typical waveguide used in modern integrated optical circuits. An example is illustrated in Fig. 2.24a). Index guided modes exist for certain frequencies of light, \( \omega \), and wavevectors, \( k_\parallel \), which are conserved. An example of the fundamental transverse electric (TE) (electric field in the horizontal plane) mode profile for silicon nanowire is shown in 2.24b).

### 2.3.2 The photonic band gap

The optical properties of dielectric media can be tailored by careful design of the sub-wavelength features. These engineered structures are called ‘photonic crystals’. Examples can be found in nature including the butterfly shown in Fig. 2.25a). The wings appear blue not because of a dye or pigment, but due to a Bragg mirror effect created by the air/dielectric nanostructure, which is periodic on the sub-wavelength scale. This structure creates a photonic band gap. Light can be confined in optical waveguides the same way, with an example shown in Fig. 2.25b). By varying the properties of the periodic structure, waveguides with unique and desirable properties can be realised, for

![Figure 2.25: a) SEM of a butterfly wing [148]. The blue colour of the wings is a result of the subwavelength structure, not of any dye or pigment. b) SEM of a silicon nanowire waveguide. The photonic crystal is 220 nm thick, with a lattice period of 500 nm and hole radius of 250 nm. One row of holes is removed to form the waveguide.](image-url)
example the slow-light effect discussed in section 2.3.3.

The idea of the photonic crystal was first introduced by E. Yablonovitch and independently by S. John in 1987 [149,150]. This section demonstrates the physical origin of the photonic band gap using a phenomenological approach. A more detailed mathematical derivation is presented in [151].

To investigate the origin of a photonic band gap, this Section begins by examining a uniform waveguide in one dimension, as shown in Fig. 2.26a). For simplicity, take the waveguide to be lossless, to be linear and to have infinite boundaries in the plane transverse to propagation. The dispersion relation for propagating in \( \vec{z} \) is

\[
\omega = \frac{c|\mathbf{k}|}{n},
\]

(2.29)

where \( c \) is the speed of light, \( |\mathbf{k}| \) is the wave vector in the \( \vec{z} \) direction and \( n \) is the index of refraction of the waveguide material. The dispersion diagram is shown in Fig. 2.26b).

Next, impose an imaginary periodicity of period, \( a \), as shown in Fig. 2.27a). All of the dispersion information can now be represented in what is called the ‘first Brillouin zone’, spanning wavevectors \(-\frac{\pi}{a} < k_z \leq \frac{\pi}{a}\). The first Brillouin zone is the smallest repeating volume in reciprocal space, the space spanned by the wavevector, \( \mathbf{k} \). The waveguide modes now fold back at the boundary of the first Brillouin zone, as seen in Fig. 2.27b).

Now, add a refractive index modulation with a period of \( a \), as shown in Fig. 2.29a). The symmetry of the structure allows only two non-degenerate modes to exist. One mode has a node centred in the region of higher refractive
index and the other in the region of lower refractive index material, because of the variation principle [151]. The first mode has more of its energy located in the higher refractive index material, which drags its frequency lower. Conversely, the other mode has relatively more of its energy in the lower refractive index material, thereby shifting the frequency higher. As a result, a gap opens at the Brillouin zone edge, where there is a degeneracy point, visible in Fig. 2.29b).

By increasing the refractive index contrast (Fig. 2.29a), the band gap opens up further. Figure 2.29b) shows the band gap for a large contrast ratio in refractive index (such as air and silicon, a ratio of 1 to 3.45). This opens a clear photonic band gap, for a large range of $\omega$.

This effect can be extended to two dimensions (2D). The first 2D optical band gap was demonstrated in 1996 by Thomas Krauss [152], who fabricated

Figure 2.27: a) 1D uniform dielectric waveguide with arbitrary period length, $a$, imposed. b) Dispersion diagram limited to first Brillouin zone, with band folding.

Figure 2.28: a) 1D waveguide with small amplitude periodic changes in the refractive index at period length $a$. b) A photonic band gap begins to open.
a hexagonal lattice of air holes in AlGaAs membrane to create a photonic band gap in the plane of the crystal. An scanning electron microscope (SEM) image of part of the photonic crystal used in that demonstration is shown in Fig. 2.30a).

The most relevant 2D photonic crystal for this thesis is a hexagonal lattice of holes in a 220 nm thick silicon membrane, with a period of \( \sim 500 \) nm and hole radius of 250 nm. In the TE polarisation, the band diagram has a photonic band gap, as shown qualitatively in Fig. 2.30c). Typically, the vertices of the reciprocal lattice vectors for the first Brillouin zone for hexagonal lattice photonic crystal are labelled \( \Gamma, K, M \) as illustrated schematically in Fig. 2.30b). Here, photonics borrows from the notation used in crystallography [153]. The X-axis in the dispersion diagram of Fig. 2.30c) is labelled with this convention.

![Figure 2.29: a) 1D waveguide with large amplitude periodic changes in the refractive index at period length a. b) A clear photonic band gap has opened with a large frequency range where propagation is forbidden.](image)

![Figure 2.30: a) An SEM image of a 2D photonic crystal [152]. The hexagonal lattice of airholes has a period of 190nm and diameter of 105nm. b) Schematic illustration of photonic crystal with a hexagonal lattice. c) Band structure for 2D photonic crystal in 220 nm thick silicon with period 500 nm.](image)
2.3.3 Slow-light photonic crystal waveguides

A waveguide can be created by placing a photonic bandgap structure on either side of a thin slab of dielectric. The result will look like a 2D photonic crystal with a row of holes removed. This was already introduced in Fig. 2.25b), where light was confined vertically by index guiding and horizontally by photonic band gap guiding. The photonic bandgap is crucial to the lateral confinement of the photonic crystal waveguide (PhCWG). For frequencies outside the bandgap light can couple to the extended modes of the photonic crystal slab and leak laterally out of the waveguide. The dispersion of a PhCWG can be finely engineered by adjusting the position of the holes in certain parts of the lattice. The spatial distribution of different modes of the waveguide are concentrated in different parts of the PhC. This spatial variation allows certain modes to be influenced more than others when engineering the structure, thereby enabling their dispersion to be somewhat independent [129,130,154].

One dramatic application of dispersion engineering is slow light [125, 155, 156]. The physical quantity of interest for a slow-light PhCWG is the group velocity. The group velocity is the speed at which a pulse envelope travels through the waveguide, given by

\[ v_g = \frac{\partial \omega}{\partial k_z}, \]

(2.30)

where \(\frac{\delta \omega}{\delta k_z}\) is the gradient of the dispersion. The ratio of the phase velocity in the bulk material, \(v_\phi\), over the group velocity, \(v_g\), gives the slowdown factor

\[ S = \frac{v_\phi}{v_g} = \frac{n_g}{n}. \]

(2.31)

Here \(n_g\) and \(n\) are the group index and the refractive index, respectively. Analogous to the refractive index \((n = c/v_\phi)\), the group index is related to the group velocity by \(n_g = c/v_g\). The higher the group index \(n_g\) of the waveguide mode, the slower a light pulse will travel through it.

Figure 2.31a) shows schematically the dispersion diagram of a silicon PhCWG, imaged in the inset. The bands are projected so that the x-axis represents the wavevector, \(\mathbf{k}\), in the direction of the waveguide. The grey region spanning
Figure 2.31: a) Schematic of a photonic crystal waveguide projected band diagram. An SEM of an example waveguide is shown in the inset. The red is the waveguide mode. Region (1) is the index-guided part of the mode, which lies outside the blue dash-dot light line. Region (2) approaches the band-edge and is highly dispersive. Region (3) is the gap-guided region, where Bragg reflection is responsible for in plane confinement. b) A close-up of area (2) for the case of no dispersion engineering. c) This region is a highly dispersive slow-light region, with the group index changing from finite to infinite at the band edge. Pulses here are chirped and broadened by the high dispersion (inset). d) A close up of area (3) for the case where a flat band is created via dispersion engineering. e) A broadband region of near-zero dispersion is created, allowing pulses to be slowed without distortion (inset).

horizontally across the middle of the figure is the projected band gap, with the bulk and air bands below and above, respectively. The red curve is the waveguide mode.

In the region marked (1), the mode is index guided and has approximately linear dispersion as would be expected in the bulk material. It reaches $k = 0$ with non-zero group velocity, $v_g$, at a Dirac like point. This region is not practically usable because the light is not confined, as it lies above the light line. The light line is the blue dash-dot line running diagonally from bottom left to top right, across Fig. 2.31a). Below this line are the regions confined for the index contrast associated with silicon and air. For more details on the light line, see reference [151].

In the region marked (2), the mode approaches the band edge. At the band edge, $v_g$ must equal zero, and $n_g \to \infty$. At $\Delta k$ away from the band edge, $n_g$
is finite, see Fig. 2.31b) and c), meaning a small change in $k$ corresponds to a large change in $n_g$. Therefore, although the light travels slowly in this region, it is highly dispersed. A pulse travelling at the frequency corresponding to the band edge becomes chirped ($\omega$ spread in $t$), as shown in Fig. 2.31c).

In the region marked (3), a guided mode with both a high $n_g$ (low $v_g$) and low dispersion can be obtained. A region of constant $\delta\omega/\delta k$ can be created by carefully engineering the position of the holes adjacent to the waveguide, (see [154] for details), as shown in Fig. 2.31d). $n_g$ is also constant here, thereby leading to a broad range of $\omega$ where $n_g$ is high and has low dispersion, see Fig. 2.31e). In this range, a pulse travelling through the waveguide is slowed without being chirped or temporally broadened.

For a near zero dispersion slow-light region (for example, see Fig. 2.31e), the Kerr nonlinearity is enhanced by two effects [156,157] – I. an increase in accumulated nonlinear phase shift – II. spatial compression leading to increased peak intensity. Effect I. occurs due to the light pulse interacting longer with the Kerr medium in the waveguide. This increased interaction time enhances nonlinear effects, similarly to having a longer waveguide. For II., the pulse is spatially compressed as it enters the slow-light region, as the front of the pulse slows down and the back catches up. To satisfy conservation of energy, neglecting loss, the compressed pulse will have an increased peak intensity further enhancing nonlinear effects.

Kerr nonlinear effects scale quadratically with the nonlinear phase shift, $\Theta$. The nonlinear phase shift $\Theta = \gamma P_0 L$, where $\gamma$ is the Kerr coefficient, $P_0$ the pump peak power, and $L$ the device length, shown in equation 2.32. The four-wave mixing efficiency, $\eta_{FWM}$, is enhanced by the effects I. and II. above, and quantitatively increases with the fourth power of the slowdown factor, S,

$$ \eta_{FWM} = S^4 \Theta^2 e^{-\alpha L}$$

where $\alpha$ the linear loss constant and $\theta$ is a phase term [130]. Having a slowdown factor on the order of ten for a non-dispersive slow-light PhCWG enhances the four-wave mixing efficiency ideally by ten thousand [129]. This allows photon pair generation by SFWM to occur efficiently in a PhCWG with a physical length of only one hundred micrometres [158]. Many of these ultra-compact
and broadband nonlinear elements can be fabricated onto a single chip offering the potential to scale-up photonic nonlinear photon generation for complex quantum photonic circuits [138].
Chapter 3

Photon pair generation in amorphous chalcogenide

This chapter considers whether it is possible to mitigate the intrinsic SpRS noise, thereby making chalcogenide a viable platform for nonlinear photon-pair generation.

This chapter is based on the following publications:


The first part of this chapter investigates the effects of temperature on photon pair generation in chalcogenide fibre. This yields limited success, highlighting the need for a new tool to better probe the SpRS in the waveguide. The second part of the chapter describes a novel tool, a broadband single-photon Raman spectrometer, used to measure the SpRS in chalcogenide, revealing the subtleties of the SpRS spectrum. This allows the region with the lowest scattering, (the low-Raman window), to be targeted for pair generation – an idea first proposed in 2010 [85]. The final part of this chapter presents photon-pair generation in a dispersion-engineered waveguide, were the photon pairs are generated in the low-Raman window, thereby enabling low noise operation. By successfully evading SpRS noise, this study demonstrates that chalcogenide remains a promising platform for future quantum applications.

3.1 Effect of cooling on Raman scattering for correlated photon-pair generation in chalcogenide

This section investigates the impact of SpRS photon pair correlations in chalcogenide glass. Unlike silica, chalcogenide is a highly nonlinear material that can be used to fabricate compact photonic chip devices [107]. Using a 7 cm As$_2$S$_3$ single-mode chalcogenide fibre a marked improvement to the CAR was attained from 0.5 at room temperature to 4.2 at a temperature of 77 K. The order of magnitude improvement at room temperature from the previous reported result [86] can be attributed to the use of a pulsed pump, while the seven times improvement is due to a reduction in SpRS noise when cooling. The CAR is characterised over a range of input powers and at different temperatures. Fitting the data with an analytical model reveals there is still a large dominant source of noise that is linearly proportional to peak power, consistent with the presence of SpRS photons. The SpRS statistics were analysed and the effect of cooling was found to be less significant as the photon channels were further detuned from the pump wavelength.
3.1.1 Fibre cooling experimental setup

The photon generation method exploited in this work is that of SFWM in a $\chi^{(3)}$ nonlinear medium. In this process, as explained in Section 2.2.2, two pump photons are annihilated to generate signal and idler photons, of higher and lower energy, respectively. These photons are correlated in time. The quality of this source is reliant on many experimental parameters that can introduce uncorrelated photons or noise into the system, such as leaked pump, detector dark counts or SpRS. The former two noise components can be avoided through filtering and improved detector technology, while the latter is a source of noise that is in the same spectral band as the signal and idler photons. The SpRS process is shown in Fig. 2.13—in this case, the Stokes process. A pump photon scatters in the medium, resulting in a lower energy photon and a phonon signified by the red arrow.

Cooling a waveguide and the required alignment stages with liquid nitrogen is technically impractical. Instead, a 7 cm long As$_2$S$_3$ single-mode fibre was used for these measurements, with chalcogenide exhibiting many desirable properties for nonlinear photon pair generation [107], including a strong Kerr nonlinearity comparable to that in silicon, but with low two-photon absorption and without the detrimental free carriers present in semiconductor materials (TPA is described in Section 2.2.4).

The fibre was fabricated by tapering a multimode fibre (CorActive), with a cladding diameter of 140 $\mu$m and a core diameter of 5 $\mu$m, through a heating and drawing process (using a soft-glass fibre tapering rig [159]) until the desired single-mode guidance could be achieved at a cladding diameter of 80 $\mu$m and core diameter of 2.8 $\mu$m. At this diameter, the nonlinear coefficient was measured to be $\gamma \sim 1.7 \text{ W}^{-1}\text{m}^{-1}$ [160] with a dispersion of $-397 \text{ ps.nm}^{-1}\text{km}^{-1}$ at 1,550 nm, and a calculated SFWM bandwidth (full width at half maximum) of 2.4 THz. This provided a nonlinear phase shift $\gamma P L = 0.059$ for a 0.5 W peak power input. The chalcogenide fibre was butt coupled directly to high numerical aperture (NA) silica fibre (refractive index of 1.44) to minimise the mode mismatch, and secured with ultraviolet-cured index matching glue (refractive index of 1.56) to reduce Fresnel reflections. The high-NA fibre was then fusion spliced to standard single-mode fibre (SMF) pigtails with FC/APC connectors.
Figure 3.1: Schematic of the experimental setup. An attenuator (ATT), tuneable band-pass filter (TBPF), and polarisation controller (PC) condition the pump pulses before coupling into the 7 cm $\text{As}_2\text{S}_3$ fibre, where correlated photon pairs are generated through SFWM. The fibre can be placed at room temperature or cooled to 77 K in a liquid nitrogen bath. The output is fed into a circulator and a FBG to block and reject the pump. An AWG separates pair photons, followed by TBPFs to further isolate pump noise. SPDs 1 and 2 detect the photons, synchronised to the pump laser by an electrical delay generator (DG). Coincidences are measured using a time interval analyser (TIA).

and placed in an insulated container to allow immersion in liquid nitrogen.

Figure 3.1 shows the experimental setup. A picosecond pulsed fibre laser generates 10 ps pump pulses at 1550.1 nm, with 0.3 nm bandwidth and repetition rate of 10 MHz. An isolator protects the laser from back reflections, while a variable attenuator and polarisation controller condition the pump pulse intensity and polarisation. A 1,550/980 nm WDM blocks any leaked 980 nm cavity pump photons. A 90/10 coupler allows the input power and spectrum to be monitored. Correlated photon pairs are generated in the $\text{As}_2\text{S}_3$ fibre via SFWM and the output is fed into a 0.5 nm bandwidth fibre Bragg grating (FBG) centred on the pump frequency to block any further pump photons. Signal and idler photons are input to an AWG separating the higher and lower
wavelength photons. AWG output channels are spaced evenly by 100 GHz, with a bandwidth of 45 GHz. A tuneable band-pass filter at each of the AWG outputs further removed leaked 1,550 nm pump photons to give a total pump suppression >100 dB. The signal and idler photons were detected by InGaAs SPDs (ID210) with a 1 ns window synchronised with the pump laser. At 20% detection efficiency, dead time of 20 µs and 5 MHz triggering rate, the detector dark count rate was ∼100 s⁻¹. SFWM gain and pump leakage both decreased as pump detuning increased, with optimum performance achieved at 700 GHz, which was used for all coincidence measurements, positioning the signal and idler at wavelengths of 1,544.5 nm and 1,555.7 nm respectively.

The CAR, defined as the ratio of correlated events to the system noise, was measured for various input pump powers. A true coincidence event involves two photons generated via SFWM in the As₂S₃ fibre and originating from the same pump pulse being detected simultaneously. This study measured accidental coincidences by looking at photons that arrived at the detectors synchronised with the pump clock but separated by one period, capturing coincidences caused by the detector dark counts, pump leakage, multiple-pair generation and SpRS noise photons. The raw coincidence histogram is plotted in Fig. 3.2a), with the coincidence, C, and accidental, A, counts corresponding to the larger and smaller peaks, respectively.

### 3.1.2 Photon statistics measurements

The results at room temperature and when the device was immersed in liquid nitrogen are shown in Fig. 3.2b). At room temperature, the maximum CAR was 0.5–an order of magnitude increase over the previously published result using a continuous wave (CW) laser in a chalcogenide waveguide [86], where the nonlinear phase shift, γPL, was similar to that used here, as the higher peak power compensated for the reduced nonlinearity of the device. The increased CAR was achieved due to the pulsed pumping providing time-of-arrival information for the generated photons.

At liquid nitrogen temperature (77 K), the maximum CAR was improved by a further order of magnitude to 4.2. Cooling affected neither the detector settings nor the coupled power; hence, the dark counts, pump leakage and
multi-pair generation remained constant. The increase in CAR can be attributed to a reduction in SpRS noise photons. Further improvements to the filtering, such as including a broadband FBG to block more pump leakage—are expected to further increase the maximum CAR. A 5 dB drop in the single-photon detection rate was observed after cooling the As$_2$S$_3$ fibre, for constant input power, which was attributed to the reduction in Raman noise.

**Photon statistics analysis**

The data were fit using a simple analytical model [123] of the CAR plots for the different temperatures to further investigate the noise. The true coincidences per pulse, $C$, in the system arose from the SFWM interaction, and can be described as $C = \eta_s \eta_i \mu$, where $\mu$ is the number of pairs generated per pulse and $\eta_{s,i}$ are the lumped collection efficiencies, including coupling efficiencies, component losses and detector efficiencies, in the signal and idler arms, respectively. Accidental coincidence counts per pulse, $A$, are defined as $A = N_s N_i$, where $N_{s,i}$ are the detected counts in the signal and idler arms. These singles counts can be written as $N_s = (\mu + \mu N_s) \eta_s + D_s$ and $N_i = (\mu + \mu N_i) \eta_i + D_i$.

![Figure 3.2:](image)

**Figure 3.2:** a) Raw coincidence count histogram. The x-axis represents the time between idler and signal-photon detector clicks. The larger peak is the coincidences, $C$, and smaller peaks are the accidentals, $A$. An artificial $\sim\!20\,\text{ns}$ delay was added to the 77 K data (blue dots) for clarity. b) CAR measured at room temperature (black squares) and at 77 K (red circles) with a signal and idler detuning of 700 GHz. Red lines represent fits to the data, as described in the text. Errors are calculated from Poissonian statistics.
where $\mu_{N_s,N_i}$ are the probability of generating a noise photon per pulse and $D_{s,i}$ are the detector dark counts, for the signal and idler arms respectively. Taking the standard formula for the $CAR = C/A$ and substituting for $C$ and $A$, the fitting function is

$$CAR = \frac{\eta_s\eta_i\mu}{[(\eta_s(\mu + \mu_{N_s}) + D_s)(\eta_i(\mu + \mu_{N_i}) + D_i)]}.$$  

(3.1)

The dark counts were measured to be $100\text{s}^{-1}$ in both detectors when triggered at the laser repetition rate of $5\text{MHz}$. The overall collection efficiencies were measured to be $3.0\%$ and $2.2\%$, for signal and idler, respectively. Using the experimental values for $D_{s,i}$ and $\eta_{s,i}$ in equation 3.1 the fitting function was used to find the free parameters, $\mu_{N_s,N_i}$. The ratio of $\mu_{N_s}$ and $\mu_{N_i}$ in the cooled and room temperature cases shows that the SpRS was reduced in both the signal and idler channels by factors of $3.16$ and $3.31$, respectively. Taking the signal and idler count rates with respect to input power in the two temperature regimes and fitting with a second-order polynomial, the ratios of the linear noise terms were $3.07$ and $3.27$, respectively, which was in good agreement with the values extracted from the CAR plots. It should be noted that although this is a significant reduction in the SpRS that resulted in a marked improvement to the CAR achieved, it is still two orders of magnitude from the ideal CAR of $\sim 315$ in the case where $\mu_{N_s,N_i}$ tend to zero.

The effect of cooling in this sample can be understood by looking at the source of SpRS—namely, the photon-phonon scattering interaction. The anti-Stokes component of SpRS is always less intense than the Stokes component; thus, the latter is the dominant source of noise affecting the correlated photon pair statistics. The number of SpRS photons in the sample is given by [161]

$$N_{SpRS}(\nu, T) = 8\pi^3 \Delta \nu P_0 L g_{DOS}(\nu)\frac{1}{\nu^2}[1 + n(\nu, T)]^\frac{1}{2},$$

(3.2)

where $\Delta \nu$ is the bandwidth of operation, $P_0$ is the peak power, $L$ is the length of the device, $\nu_{Stokes}$ is the frequency of the SpRS photons, $\nu$ is the detuning from the pump and $T$ is the temperature. $g_{DOS}(\nu)$ is the phonon density of states, shown in Fig. 3.3a), which is directly proportional to the Raman gain, with $g_{Gain}(\nu) = \frac{1}{\nu}g_{DOS}(\nu)$. The phonon population, $n(\nu, T)$, is described by
the Bose-Einstein distribution

\[ n_{\text{BE}}(\nu, T) = \frac{1}{e^{\frac{\hbar \nu}{k_B T}} - 1}. \]  

(3.3)

The photon population distributions are shown in Fig. 3.3b) for the room temperature (293 K) and liquid nitrogen cooled cases (77 K). For small detunings the phonon occupation probability is high meaning the lower energy states are more likely to be occupied. A high occupation probability means that, despite a low density of states, a large number of photons are scattered in the small detuning region. This is seen in Fig. 3.3c), where the SpRS intensity increases close to the pump.

In the far detuned cases, \( \nu > 6 \text{THz} \), the phonon population tends towards zero, and the high and low temperature cases converge. Here cooling the device has a negligible effect on the SpRS intensity. It should be noted that there is a low-Raman window at 7.4 THz detuned from the pump–shown in purple in Fig. 3.3c)–where cooling would not be required in a device to exhibit a high CAR.

The effect of Raman scattering on photon-pair generation in chalcogenide was investigated by measuring photon statistics at different temperatures. Cooling the device in liquid nitrogen, at a pump detuning of 700 GHz, in-

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**Figure 3.3:** a) The Raman density of states for chalcogenide, \( g_{\text{DOS}} \propto \frac{1}{\nu} g_{\text{gain}}(\nu) \), extracted from the spectrum of a bulk As\(_2\)S\(_3\) glass sample measured using the 90° scattering method in reference [162]. b) The Bose-Einstein phonon population distribution at 293 K and 77 K. c) Lower energy states are more likely to be occupied, thereby enhancing the SpRS close to the pump. The low-Raman window–an intrinsic minimum for Raman scattering in chalcogenide–is highlighted in purple.
creased the signal-to-noise ratio (CAR) by an order of magnitude. By fitting an analytical model to the experimental data it was shown that the SpRS photons were reduced by approximately three times. Although the cooling made a marked difference in the near-detuned regime, it did not increase the CAR to the desired level of over 10, and was far from the ideal value of approximately 315. In the far detuned region cooling was shown to have a negligible effect on the number of Stokes SpRS photons, and will thus provide little improvement.

In the region of low SpRS at a photon-channel detuning of around 7.4 THz, improved photon-pair statistics can be expected without the requirement of cryogenic cooling. However, accessing this window requires the SFWM bandwidth be extended. Photon correlation measurements are an indirect observation of the SpRS. Before targeting the low-Raman window for pair generation, a direct measured of the SpRS and the effect of cooling, is required. The next section presents the direct measurement of the SpRS spectrum.
3.2 Broadband single photon Raman spectroscopy

The previous section yielded a modest improvement to a chalcogenide based photon pair source through cooling. This highlighted the need for a new tool for characterising the SpRS to enable more intelligent mitigation of this intrinsic noise. Thus, this section presents a novel method for directly measuring the SpRS in optical waveguides, in an alignment-free setup. The broadband photon-counting Raman spectrometer consists of a pulsed laser, liquid-crystal-on-silicon (LCoS) spatial light modulator used as a reconfigurable optical filter and a SPD, sketched in Fig. 3.4. The temperature and polarisation dependence are characterised in an As$_2$S$_3$ amorphous glass fibre for a Stokes detuning range of 1 to 9 THz from the pump frequency. The experimental data were fit with a theoretical model and this fit was used to extract the Raman-gain spectrum, which was compared with measurements of bulk As$_2$S$_3$. The sensitivity of the method allows direct characterisation of chip-scale photonic devices. Knowledge of the SpRS spectrum enables the choice of lowest noise configuration for the pump and photon channels [163]; however at the single-photon level, SpRS cannot be detected using a standard spectrum analyser.

Spontaneous and stimulated Raman scattering were measured in order to characterise both Raman noise and material gain spectra. The spectral distribution of SpRS can be measured in bulk samples using the 90° scattering method [161, 162, 164, 165], where a free-space beam is incident on the sample and the SpRS component that is emitted perpendicular to the pump is collected by free-space optics and measured with a spectrograph.

![Figure 3.4: Method for broadband photon-counting Raman spectroscopy in optical waveguides.](image)
can also be measured in optical fibres; however, this requires a pulsed laser and long fibre lengths to achieve a measurable signal strength [166–168]. Direct measurements of stimulated Raman scattering have been performed using nonlinear pump-probe techniques [160,169–172]. All techniques described require high powers incompatible with the damage threshold of many on-chip devices and, in the nonlinear case, the addition of either a highly-tuneable or ultra-broad bandwidth probe. Recently, photon-counting techniques have been demonstrated to measure weak SpRS signals for very small detunings and low temperatures in silica fibres [98,173]. However, no method applicable to measuring SpRS directly over a broad bandwidth in chip-scale waveguides has been presented.

3.2.1 The spectrometer

The method outlined in Fig. 3.4 uses a pulsed laser to excite SpRS in the waveguide, a TBPF resolve the spectral distribution, and a SPD to provide the sensitivity needed to measure short devices. This technique was used to measure the temperature and polarisation dependence of the SpRS spectrum in a chalcogenide glass fibre [107]. By fitting the experimental data, the Raman-gain spectrum could be extracted and was found to be in good agreement with previous SpRS measurements [161]. In principle the single-photon sensitivity allows SpRS to be measured directly in chip-scale photonic devices, where 90° scattering and nonlinear techniques are not applicable.

Figure 3.5 shows the full experimental schematic. As a first demonstration, a short fibre was used for ease of immersion in cryogenic liquid; however, this technique is equally applicable to planar waveguides. The laser generated 10 ps pulses centred at 1,535 nm, with a 0.3 nm bandwidth at a repetition rate of 10 MHz. An isolator protected the laser from back reflections and a variable ATT with a PC conditioned the pulses. Two cascaded 1,550/980 nm wavelength division multiplexer modules blocked any leaked 980 nm cavity pump photons. An inline polariser fixed the input polarisation to the fibre. A TBPF centred at the pump was placed directly before the As$_2$S$_3$ fibre to suppress any Raman photons scattered in the preceding silica fibre.

The single-mode As$_2$S$_3$ fibre was 20 cm long with a fibre nonlinearity pa-
parameter of $\gamma = 1.7 \, \text{W}^{-1} \text{m}^{-1}$, refractive index of 2.44 and transmission loss of 1 dB m$^{-1}$. The As$_2$S$_3$ fibre was housed in an insulated container that was filled with liquid nitrogen for the 77 K measurements. At the As$_2$S$_3$ fibre output, a PC and PBS were used to select either the parallel or perpendicular polarisation component.

Previously, the measurement bandwidth was limited by the challenge of creating a widely tuneable passband with more than 50 dB isolation [98, 173]. A 100 GHz wide flexible pass-band with 90 dB of out-of-band noise suppression was established using a LCoS spatial light modulator (Finisar Waveshaper) [174] in combination with a tuneable C/L band BPF. The SpRS spectrum was mapped by moving the filter pass-band to different detunings, and counting the generated photons. The component loss varied over the full detuning range, and was corrected for at each detuning. The Waveshaper was characterised using an amplified spontaneous emission (ASE) source before it was inserted into the setup. To measure the BPF, a 99%/1% coupler was placed between the Waveshaper and TBPF. A C/L band source of broad-band ASE was connected to the 1% input and the TBPF at the 99% output. A second coupler was
connected after the BPF, with Output 1 directed to a SPD (Id-Quantique ID210) via an ATT, and output 2 connected to an optical spectrum analyser. By blocking the SPD and opening the ASE source, the C+L BPF was manually tuned and the loss (including couplers) was measured for each detuning without breaking any fibre connections. The quantum efficiency of the detector at each detuning was also taken into account. From this, a frequency-dependent transfer function was calculated for the experimental setup.

The probability of scattering a photon from the pump to the Stokes field, within the measurement bandwidth of \(\Delta\nu = 100\ \text{GHz}\) was less than 0.01 SpRS photons per pump pulse, inferred from the photon count rate; therefore, it was appropriate to count photons using an SPD. This corresponded to a coupled peak power at the fibre of 1.6 W, assuming symmetric input-output coupling loss. The detectors were triggered in synchrony with the pulsed pump at 10 MHz and set to a detection efficiency of 10\%, effective gate width of 1 ns, and dead time of 20 \(\mu\)s. The SPD characteristic dark count rate was approximately 50 Hz, when triggered at 10 MHz. To check for pump leakage, an ATT with comparable length pigtails was temporarily inserted into the setup in place of the As\(_2\)S\(_3\) fibre. The attenuation was set to match the coupling loss, and no increase in the noise count rate was seen over the intrinsic detector dark counts, proving pump leakage to be negligible.

As seen in Section 3, SpRS occurs when pump photons interact with the material through the polarisability of the molecular lattice and emit/absorb a phonon that lowers/increases the photon energy in the Stokes/anti-Stokes case. The experiment operated in the regime of pure SpRS, as there were initially no photons in the Stokes mode to induce the Raman scattering process. Another consideration was photons generated by SFWM in the fibre for both Stokes and anti-Stokes detuning. The dispersion parameter \(D\) was calculated to be \(-405.8\ \text{ps km}^{-1}\text{nm}^{-1}\) at 1,535 nm, corresponding to a calculated SFWM bandwidth of 0.8 THz. The SpRS measurements began at 1 THz; thus, the SFWM contribution can be considered small up to 2 THz and negligible beyond that detuning [85]. Note that the bandwidth is defined to be the half-width at half maximum.
3.2.2 Single photon Raman spectrum

The measurements for the component polarised parallel to the pump at 293 K and 77 K are plotted in Fig. 3.6, as red circles and blue triangles, respectively. For near detunings, cooling the fibre to 77 K reduces the SpRS by more than 4 dB. As the frequency detuning from the pump is increased, changes in temperature, $T$, have less of an effect, with the SpRS rate for both temperatures converging. These effects can be explained by the dependence of the SpRS rate on the phonon population. SpRS can be described as the emission or absorption of a phonon by a pump photon. The Stokes case includes both spontaneous and stimulated emissions of a phonon, where the stimulated component depends on the phonon number, with a distribution described by Bose-Einstein statistics as

$$n_{BE} \propto \frac{1}{e^{h\nu/k_BT} - 1}, \tag{3.4}$$

with $\nu$ being the detuning from the pump frequency. This is not to be confused with the stimulated Raman scattering process enhanced by the presence of photons in the scattered mode. Thus, the Stokes intensity is proportional to $(n_{BE} + 1)$. From equation 3.4, it follows that the low energy phonon states are more likely to be occupied, enhancing the photon count rate closer to the pump, as observed in Fig. 3.6 for both temperatures. The SpRS component proportional to $n_{BE}$ is dependent on $T$ and is strongest close to the pump. This dependence of the SpRS rate on $T$ weakens for increased detuning as the ratio between the phonon energy, $h\nu$, and thermal energy, $k_BT$, increases and $n_{BE}$ becomes less dominant. Note that, even at zero temperature Stokes scattering occurs via spontaneous phonon emission.

Polarisation effects

Polarisation effects are observable in the Raman scattering spectrum even in the absence of material anisotropy. In an isotropic medium, there are isotropic and anisotropic components corresponding to the two free parameters of the Raman response tensor [163,175]. Figure 3.7 shows the SpRS component polarised parallel (green diamonds) and perpendicular (brown squares) to the
Figure 3.6: Measurement of the SpRS photon count rate with respect to detuning from the pump frequency at room temperature (293 K, red circles) and at liquid nitrogen temperature (77 K, blue triangles).

pump, measured at room temperature. These orthogonal polarisation components include different contributions from the isotropic and anisotropic Raman response, and thus have different shaped spectra, as previously observed for As$_2$S$_3$ glass [161]. The SpRS photon count rate [163] is expressed as

$$R_{\text{SpRS}}(\nu, T) = C \eta(\nu_s) \Delta \nu P_0 L \left[ 1 + n_{\text{BE}} \right] g_{\text{gain}}(\nu),$$  \hspace{1cm} (3.5)

where $C$ is the Raman coupling coefficient; $\eta(\nu_s)$ is a frequency-dependent transfer function for the experimental setup; $\nu$ the detuning from the pump; $\Delta \nu$ is the fixed measurement bandwidth; $P_0$ the coupled peak power; $L$ the effective fibre length; $\nu_s$ the Stokes frequency; and $g_{\text{gain}}(\nu)$ the Raman-gain spectrum, which is determined by the Fourier transform of the Raman response functions [176]. The Raman-gain includes the polarisation components parallel and perpendicular to the pump $g_{\text{gain}}(\nu) = g_{\parallel} + g_{\perp}$. By fitting the measurement using equation 3.5 the Raman-gain, $g_{\text{gain}}(\nu)$, was extracted and compared to data from measurements of a bulk sample of As$_2$S$_3$ glass. The normalised gain for the parallel polarisation, $g_{\parallel}$, extracted from the SpRS measurement and the gain measured in a bulk sample are shown in Fig. 3.8 as red circles and a black dashed line, respectively. The two spectra are in good agreement, which further confirms the accuracy of the method. The measurement was limited by the C/L band TBPF and thus could not extend beyond
9 THz from the pump. For small detunings the Bose-Einstein equation 3.4 will increase approximately proportional to $\frac{1}{\nu}$. It has previously been observed that the SpRS rate does not follow the Bose-Einstein distribution to infinity, but instead goes to zero at the pump [91,161,162]. Therefore, to lowest order, the gain (which is an odd function of $\nu$) must vary as $\nu^3$ close to the pump. The SpRS was not measured closer to the pump because the signal was contaminated with SFWM and pump leakage. The extracted Raman-gain, $g_{\text{gain}}(\nu)$, can be used to calculate the anti-Stokes SpRS for the As$_2$S$_3$ fibre, as shown in Fig. 3.9. The existence of the low-Raman window at 7.4 THz detuning was confirmed for both the Stokes and anti-Stokes emission. The window is the result of a characteristic dip in the phonon density of states [161].

**Discussion**

Possible improvements to this method include increasing the measurement range of 9 THz, which is currently limited by the C/Lband TBPF and C-band pump. This could be implemented using a tunable source, such as an optical parametric oscillator, to pump at a shorter wavelength and extend the maximum detuning range. For integrated devices where broad SFWM would contaminate the SpRS signal, the back scattered SpRS could be measured and used to predict the forward Stokes and anti-Stokes SpRS noise. The SFWM contribution could also be identified in the forward scattering by including a
second detector at the anti-Stokes wavelength to measure coincidence counts [91], with the coincident photons originating from the four-wave mixing.

The ability to characterise the Raman scattering is useful in quantum communication applications, where, for example, the in-band noise of a single-photon channel operating near a bright laser can be reduced by moving it to a known SpRS minimum [177]. This could apply to photon-pair generation via SFWM, frequency conversion via Bragg-scattering four-wave-mixing [178] (as discussed

Figure 3.8: The Raman-gain $g_{\text{gain}}(\nu)$ (red circles) for As$_2$S$_3$ glass extracted from the SpRS measurements and superimposed onto a spectrum from bulk As$_2$S$_3$ measured using the 90° scattering method (dashed black line) [162].

Figure 3.9: The measured Stokes and predicted anti-Stokes SpRS for room temperature (red circles) and 77K (blue triangles). The anti-Stokes calculation used the extract Raman-gain for the parallel polarised component of the Stokes SpRS.
in more detail in the section 5.2) or a classical information signal transmitted adjacent to the quantum channel [179].

3.2.3 Raman in silicon

Another material of interest is CMOS-compatible amorphous silicon (which is free from two-photon-absorption) and its crystalline counterpart. Silicon-integrated photonic devices can have characteristic sizes ranging from a few millimetres to hundreds of nanometres. Typical techniques for measuring the spectral distribution of SpRS in bulk samples with the free-space 90° scattering method [180] are applicable to measuring these devices. More importantly, for a hybrid structure that includes additional materials onto a silicon core device, Raman spectroscopy could be a vital characterisation tool. In long fibres, the SpRS spectra can be measured using a pulsed laser to achieve measurable signals [166]. High power pulses are likely to have an energy density incompatible with the damage threshold of many on-chip devices. Measurements of stimulated Raman scattering have been performed using nonlinear pump-probe techniques [181], requiring the addition of either a highly tuneable or ultra-broad bandwidth probe. Recently, photon-counting techniques have been demonstrated to measure weak SpRS signals in fibres [98, 173, 182]; however, no direct measurements of the SpRS spectra of nanophotonic-chip devices over a broad bandwidth have been performed.

Using a variation to the setup from Fig. 3.5, the SpRS spectrum was measured for both amorphous silicon (a-Si) and crystalline silicon (c-Si) nanowire devices, with both platforms being CMOS-compatible. The method uses a CW laser to excite SpRS in the waveguide, a spatial light modulator (Finisar waveshaper) to resolve the spectral distribution, and a SPD to provide the sensitivity needed to measure the small signal associated with the SpRS in chip-scale waveguides. The results shown in Fig. 3.10 display a narrow Raman peak for c-Si at the frequency shift of $\sim 525 \text{ cm}^{-1}$ with a bandwidth of $<150 \text{ GHz}$, with typical reported values of $\sim 100 \text{ GHz}$.

In contrast, the Raman spectrum for a-Si has a 3 dB bandwidth of $\sim 2.5 \text{ THz}$, comparable to the bandwidth of $\sim 6 \text{ THz}$ for silica glass the standard material used in telecom distributed Raman amplifiers. As the SpRS spectrum is pro-
portional to the Raman gain [182], the a-Si platform proves a promising candidate for broadband CMOS compatible Raman amplifiers. In addition, a-Si displays low TPA [183] and free carrier generation—the main limiting factors in the efficiency of c-Si Raman amplifiers [184].

Figure 3.10: The normalised SpRS spectrum for amorphous and crystalline silicon nanowires measured on the polarisation axis parallel to the pump, in spectroscopy units (cm$^{-1}$).
3.3 Photon generation in the low-Raman window

Having measured the SpRS using the novel technique from the previous section, this section demonstrates correlated photon-pair generation at the low-Raman window in a dispersion-engineered 10 mm As$_2$S$_3$ chalcogenide waveguide. This is a dramatic improvement of those previously achieved [86], with a room temperature CAR of 16.8—more than two orders of magnitude improvement on previous on chip measurements [86]. Building on previous theoretical work [85], this result was made possible by carefully engineering the waveguide dispersion for broadband SFWM to allow the photon pass-bands to be placed in the low-Raman window [161].

![Figure 3.11: a) Schematic of dispersion engineering goal of moving the photon pair generation into the low-Raman window. b) Image of the 10 mm long dispersion engineered chalcogenide waveguide coupled to lensed fibres.](image)

3.3.1 Dispersion engineered chalcogenide waveguide

The device, illustrated in Fig. 3.13 insert, is a 10 mm long rib waveguide, that is 2 $\mu$m wide and 350 nm high above a 500 nm As$_2$S$_3$ layer, deposited on a silica-on-silicon substrate and overclad with a polymer layer. The propagation loss is 0.3 and 0.8 dB.cm$^{-1}$ in the TE and transverse magnetic (TM) modes, respectively, with 5.0 dB per facet insertion loss. The nonlinear coefficient, $\gamma$, was measured to be 10 W$^{-1}$m$^{-1}$, and the dispersion, $\beta_2$, calculated at the pump wavelength of 1,545.36 nm, was $-209$ ps.nm$^{-1}$km$^{-1}$ and 25.5 ps.nm$^{-1}$km$^{-1}$ for the TE and TM modes. In longer devices the propagation loss of the waveguide restricted experiments to TE operation, limiting the SFWM bandwidth to the
region $\pm 1.4$ THz about the pump where the SpRS efficiency was high [86].

The key here is the tailoring of the dispersion of the waveguide and the use of a shorter device, thereby enabling TM mode operation and SFWM over a bandwidth of up to 9 THz. The TE and TM SFWM pair generation rate

$$f_{\text{SFWM}} = \Delta \nu (\gamma P_0 L)^2 \text{sinc}^2 \left( \frac{\beta_2 (2\pi \nu)^2}{2} L + \gamma P_0 \right), \quad (3.6)$$

is shown in Fig. 3.12, where the characteristic filter bandwidth is $\Delta \nu$, detuning from the pump is $\nu$, the group velocity dispersion is $\beta_2$ at the pump wavelength, the device length is $L$ and the peak pump power is $P_0$. Equation 3.6 assumes the pump is sufficiently weak that multi-pair events are negligible.

Now the SpRS photon generation rate [161]

$$f_{\text{SpRS}} = \Delta \nu P_0 L g(\nu) [1 + n(\nu, T)]^{1/\nu}, \quad (3.7)$$

is also shown in Fig. 3.12 and includes the pump frequency, $\nu_p$; the phonon
density of states, \( g(\nu) \); temperature, \( T \); and the boson occupation statistics,

\[
n(\nu, T) = \frac{1}{e^{\frac{\hbar \nu}{k_B T}} - 1}.
\]

(3.8)

The minimum in the SpRS spectrum results from a characteristic dip in the optical phonon density of states for amorphous \( \text{As}_2\text{S}_3 \) glass. This low-Raman window, together with the broadband SFWM enabled by the dispersion engineering, allows the efficient generation of low-noise correlated photon pairs in the TM mode [85], as shown in the shaded region of Fig. 3.12.

### 3.3.2 Photonic statistics measurements

Figure 3.13 shows the experimental setup. A mode-locked tuneable fibre laser (Pritel) centred at 1,545 nm produced 10 ps pulses at a repetition rate of 50 MHz, with peak power adjustable between 0.1 and 1 W. These passed through an isolator, followed by a 1,550/980 nm wavelength division multiplexer to block any residual cavity pump photons. An ATT and PC conditioned the pulses before reaching a BPF centred at 1,545 nm, positioned directly before the chip to block any SpRS photons generated in the preceding silica fibres. The pump pulses were then coupled into the waveguide using a lensed fibre. A coarse wavelength division multiplexer separated the signal

![Figure 3.13: Experimental setup photon-pair generation in a As2S3 waveguide. CWDM = coarse wavelength division multiplexer. IPG = Inorganic polymer glass. Insert: Cross section of As2S3 waveguide.](image)
and idler, providing 50 dB of pump suppression. A four-port circulator with two narrow-band apodised point-by-point FBGs used in reflection [185] filtered the signal photons at 1,489 nm, while a TBPF was used for the idler photons, at 1,605 nm, providing an additional 50 dB of out-of-band noise suppression. Signal and idler photons were detected by InGaAs SPDs (ID210) with a 1 ns detection window synchronous with the pump laser. At 7.5% and 10% detection efficiency, dead time of 10 µs and 50 MHz triggering rate, the detector dark count rates were $\sim 40 \text{s}^{-1}$ and 120 s$^{-1}$ for signal and idler, respectively. Different efficiencies were used to balance to total collection efficiency, unbalanced by differences in filter insertion loss, between the idler and signal arm.

**Coincidence measurements**

The CAR is again defined as the ratio of the rate of correlated event counts, $C_{true}$, to the rate of noise counts, $A$. A high CAR is desirable, indicating a better signal-to-noise ratio. A coincidence event occurs when two photons are detected with timing correlated to the same pump pulse, plotted in Fig. 3.14

![Figure 3.14: a) Coincidence rate versus peak power in the low Raman window at $\nu = 7.4 \text{THz}$. Here, the true coincidence count rate, $C$, approaches the total coincidence rate, $C_{raw}$, as SFWM is dominant over noise, $A$. b) For small detuning, $\nu = 1.9 \text{THz}$, $C$ and $A$ are comparable, as SpRS is significantly more efficient at this detuning. Poissonian error bars are equal to the size of markers.](image-url)
as triangles. Detection events separated by an additional pump period are accidents—shown in Fig. 3.14 as squares—and include detector dark counts, pump leakage, multiple-pair generation and SpRS noise photons. The true coincidence rate, $C_{\text{true}}$, is the rate of detection for photon pairs generated by SFWM and given by $C_{\text{true}} = C - A$, where $C$ is the total detected coincidence rate. In Fig. 3.14a) $A$ is negligible and $C_{\text{true}}$ approaches $C$, as SFWM is dominant when operating in the low Raman window. In contrast, in Fig. 3.14b) $A$ is comparable to $C$ as SpRS introduces uncorrelated noise when operating close to the pump.

Figure 3.15 plots the CAR as a function of the input power. The signal and idler channels were positioned in the low-Raman window at 7.4 THz either side of the pump, plotted as blue circles. A maximum CAR of 16.8 was reached, with a pair generation rate of $10^5$ per second in the waveguide and nonlinear phase shift of $\gamma P_0 L = 0.05$. Further improvement to the CAR was limited due to loss introduced by multiple passes through the four-port circulator and poor detection efficiency at the longer wavelength. Shifting the photon passbands to 1.9 THz from the pump and using appropriate filters, the maximum CAR was 0.7, shown in Fig. 3.15 as red diamonds. However, the coincidence count rate increases due to a slight improvement in both the overall collection and idler channel detection efficiencies.

![Figure 3.15: Measured CAR for small pump detuning (red diamonds) and operating in the low-Raman window (blue circles) with Poissonian errors bars.](image)

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It should be noted that the CAR of 0.7 measured for small detuning is still an order of magnitude improvement from the previous CAR of 0.07 measured in a CW experiment at similar detuning [86], due to the ability in this experiment to gate the detectors in synchrony with the arrival of the generated photons. The additional 25 times improvement is due to the dispersion engineering allowing access to the low Raman window.

\(g^{(2)}(0)\) correlation measurement

To measure the second-order correlation function, \(g^{(2)}(0)\) [9], an optical fibre 50/50 coupler was included in the idler detection channel and replaced the single ID210 detector with a pair of ID201 detectors at both outputs of the coupler. The efficiency of the signal detector was set to 25% with a dark count rate of 3.5 kHz, while both idler detectors were set to 20%. The dark count rates, at a heralding rate of 6.5 kHz, were 2.6 s\(^{-1}\) and 3.6 s\(^{-1}\) for the idler detectors. For an input power corresponding to 0.005 pairs per pulse, the second order correlation \(g^{(2)}(0) = 0.25 \pm 0.13\). This was well below the criterion of \(g^{(2)}(0) = 0.5\) required to demonstrate single photon anti-bunching. Below this power, the dark counts became dominant in the \(g^{(2)}(0)\) measurement. This confirmed that the source was operating in the quantum regime and that it could be used as a heralded single-photon source.

Discussion

This chapter has shown that it is possible to mitigate the intrinsic SpRS noise in a chalcogenide photon pair source. Cooling was shown to provide only a moderate reduction in SpRS close to the pump, and a negligible reduction far from the pump. A novel single-photon Raman spectrometer was then used to characterise the SpRS in chalcogenide and other materials. This led to low Raman-noise correlated photon pairs being generated in the low Raman window using a 10 mm long dispersion engineered planar waveguide. A CAR of 16.8 was measured at room temperature using a pulsed pump, which was a 250 times improvement of the original CW pumped As\(_2\)S\(_3\) waveguide experiment using small detuning. This work shows the potential for As\(_2\)S\(_3\) waveguides to be used as a platform for quantum photonics and communications. These results
not only apply to the generation of correlated photon pairs, but also to photonic manipulation through the use of As$_2$S$_3$ chalcogenide glass to perform quantum state translation [89] and coherent photon conversion [186] where the single photon and pump channels must be carefully placed to avoid contamination by SpRS.
Chapter 4

Spatial multiplexing of silicon nonlinear photon sources

This chapter asks whether it is possible to combine heralded single-photons, generated from monolithic devices, using active spatial multiplexing.

This chapter is based on the following publications:


This chapter demonstrates the first implementation using integrated nonlinear single-photon sources. It presents an enhancement in the multiplexed single-photon output for both dual and single-input silicon photonic crystal waveguide sources. The setups are all fibre coupled, stable and robust implementations of spatial multiplexing, taking the first step towards an on-demand photon source. Further, Section 4.3 uses an analytical model of the photon statistics to determine the multiplexing enhancement. Section 4.3.2 uses a matrix method analysis to investigate the scalability to many sources.

4.1 Improving the statistics of nonlinear photon sources

It has been demonstrated that the quantum nature of single photons could be used to create a quantum computer [44] and to communicate securely [187]. These initial proposals were formulated with the assumption that perfect photon sources would one day exist. Imperfect non-deterministic single-photon sources have limited the complexity of demonstrations to date, despite significant work to simplify quantum operations and ease technical demands [188]. If the single-photon output for multiple indistinguishable sources can be engineered to be more deterministic, this will create a new potential for scalable quantum technology based on photon–photon interference, including boson-sampling processors [189–192], long-distance communication [193] and metrology [194]. Several physical systems can be used to generate single photons. Single quantum emitters such as quantum dots [60, 61] or diamond colour centres [63] potentially offer on-demand single photons, but are difficult to implement on-chip and scale up, owing to challenges in engineering such sources to emit at the same frequency [61], the large-footprint of bulk photon collection apparatuses and the typical requirement for cryogenics [60, 61].

Attenuated coherent light sources are frequently used as a source of non-deterministic single photons [52,193]. They are easy to construct and package, but the output signal-to-noise is fundamentally limited by the Poisson shot noise of a coherent source; thus, a fraction of events will contain more than one photon and the probability of such multiple photons increases with the emission rate. One must inevitably accept either a low photon rate or a significant fraction of multiple photons. A third class—heralded single-photon sources—are
based on correlated photon-pair generation. By detecting one photon from a
pair, the existence and timing of the remaining photon is then known. This
additional information is the basis of multiplexing [132], where heralded single
photons from multiple sources can be deterministically combined using active
switching.

4.1.1 Multiplexing of heralded single-photon sources

Heralded single-photon sources are often based on photon pairs generated in
nonlinear optical media, including SFWM in $\chi^{(3)}$ waveguides [83,121,126,131,
177,195,196] and SPDC in $\chi^{(2)}$ devices [197,198]. As stochastic processes,
these schemes do display multi-pair events. If a sufficiently narrow output fil-
ter ensures that the generated photons are spectrally indistinguishable, which
is usually desirable, the number of photon pairs generated per pulse is gov-
erned by thermal photon statistics [64,189,199,200]. For weaker filtering, the
statistics approach a Poisson limit, as multiple pairs are most likely to be
produced in uncorrelated events. Either way, multi-pair generation introduces
noise photons, thus creating a fundamental relationship between the number
of photon pairs generated per pulse and the noise performance of the source,
discussed in detail in Section 2.2.2.

Now, let $\eta$ be the collection efficiency or probability a photon goes through
the circuit to the detector. For an integrated device working in the telecommu-
nications wavelengths, $\eta$ is typically between $10^{-2}$ and $10^{-3}$ because of coupling
loss, component loss and imperfect detection. The probability that one pho-
ton of a pair is detected per pulse is then $\mu \eta$, meaning that, for detecting a
heralded single-photon or two-photon coincidence per pulse, the probability
becomes $\mu \eta^2$. Next, the probability of detecting interference events between
the heralded photons from two separate sources is $\mu^2 \eta^4$, meaning that scalable
multi-photon operations are extremely uncommon. As this interference be-
tween photons from separate sources is at the core of many quantum gates, the
low probability is detrimental to the implementation of many algorithms [201].

Figure 4.1 presents a schematic of this study’s proposed architecture for
efficient spatial multiplexing of photon pair sources, with the aim of building a
more deterministic single-photon source. Here, the pump pulses are coupled to
silicon waveguides, where photon pairs are generated by SFWM. The PhCWGs provide strong light confinement, but the major enhancement to the nonlinearity is through slow-light engineering [154], in which the precise arrangement of the photonic crystal holes induces very strong dispersion, so that the pump group velocity, \( c/n_g \), is lower than the value \( c/n_0 \) in unpatterned silicon by a factor of \( S = n_g/n_0 \sim 10 \). The SFWM process scales as an \( S^4 \) enhancement owing to the increase in pump intensity and interaction time, thereby enabling an ultra-compact source [130]. To permit the generation of indistinguishable photon pairs using a single pump laser, the position of the slow-light frequency window must be nominally identical in all PhCWGs (see Fig. 4.5b).

The two photons are separated using wavelength division components that can be realised using, for example, integrated arrayed waveguide gratings [202, 203]. One photon from each pair is detected by an array of \( N \) integrated silicon waveguide-based SPDs [204], while the other is buffered in an optical delay line (ODL) [205, 206]. Currently, this delay line is implemented using

\[ \text{Figure 4.1: Scheme for spatial multiplexing of } N \text{ photonic crystal waveguides. Inset: This work’s first implementation used a device with two separate, monolithically fabricated PhCWGs, designated A1 and B1} \]
optical fibre, because this is the only technology available to provide sufficient delay, while remaining at a reasonably low loss. While the photon is delayed, a radio frequency logic gate triggered by a photon detection feeds forward a drive signal to a fibre-coupled opto-ceramic switch, which actively routes single photons to a common output. The switch, which is based on an electro-optic polarisation rotation [207], is fabricated from lead lanthanum zirconium titanate (PLZT) and has a maximum switching rate of 1 MHz, limited by the particular electronic driver, and an optical insertion loss of $\sim 1\, \text{dB}$. For a net increase in the heralded photon rate, the total switch loss must be $< 3\, \text{dB}$ (see [132]). The output is a multiplexed stream of indistinguishable single photons, assuming matched post-generation spectral filtering.

Multiplexing many photon pair sources decouples the relationship between the desirable single-photon output and multi-pair noise governing an individual source, thereby allowing a higher pair rate for a fixed multi-pair (noise) rate. Multiplexing can be implemented with both spatial [132–134] or temporal schemes [142]. The previous demonstration of efficient multiplexing used a free-space scheme [134], where the use of bulk MZIs made the setup inherently sensitive to alignment stability, and used a large amount of space, which caused the scalability of this scheme to be an issue. In contrast, the advantages of an integrated photonic architecture are that many sources and components can be fabricated onto a single monolithic chip, with proven stability [4]. Although this strategy has been highly successful for tuneable quantum processing [5], before the study presented here, there had been no demonstration for scalable multiplexed single-photon generation.

### 4.1.2 Experimental setup

In the first experiment, 7 ps laser pulses were separated off-chip with a fibre-based directional coupler, and coupled to on-chip polymer access waveguides using a lensed fibre. For the second experiment, a Y-split coupler was integrated on to the same chip as the PhCWGs, and the pump was coupled to the device via a single lensed fibre and input polymer waveguide. Note that in both experiments the pump laser repetition rate was 50 MHz. The pulses were thus divided and then coupled into the two PhCWGs, which had
nominally identical dispersion properties and slow-down factors, \( S \sim 10 \), as shown in the inset of Fig. 4.2. In the PhCWG regions, correlated photon pairs were generated via SFWM, whereby two photons from the pump pulse were annihilated to generate signal and idler photons of higher and lower energy, respectively, which must obey both energy and phase-matching conditions. The efficiency of generating photon pairs was enhanced by the slow-light effects in the PhCWGs [86], where there is a simultaneous enhancement to the peak power in the device as the light slows down, and a longer interaction time as the slowed pulse propagates through the nonlinear medium. Care must be taken not to operate in a regime of enhanced nonlinear loss [70]. Each pair was coupled out of the device, again using inverse tapers and lensed fibres to minimise loss. The photons from each pair were separated into two SMFs using a pair of arrayed waveguide gratings with a photon-channel detuning of 300 GHz from the pump channel. The signal photons were detected using superconducting-SPDs (SSPDs). The SSPDs, with a maximum allowed count rate of 10 MHz, were operated at 10% efficiency with dark counts of 100 s\(^{-1}\) and integration time ranging from five to fifteen minutes for the highest and lowest powers, respectively. These detection events provided the heralding electrical pulses that triggered the optical switch. A ceramic optical switch, made from ultra-low-

![Figure 4.2: Spatial multiplexing scheme. A single laser is coupled to a monolithic silicon chip to pump an ultra-compact array of silicon PhCWGs. Photon pairs are generated in the PhCWG region, with the wavelength separation implemented using integrated AWGs and the heralding photons detected using SPDs. The remaining photons go through a delay line, while a fast electronic logic gate sets the state of the PLZT switch. The selected heralded photon is then routed to the common fibre output.](image)
loss PLZT [207] was used to route photons from two monolithic silicon sources to a common output. Depending on whether source A or B generated a pair, the optically transparent ceramic switch routed the remaining idler photon from each pair to the common output fibre, where they were detected using a third SSPD. All counts were then analysed for coincident detection using a TIA, and measured as a histogram (see Fig. 4.10 for tabulated results).

Coincidence counts were taken in a histogram with respect to the delay time, as described in Section 2.2.2. Coincidences at zero time delay included the desired contribution from single pairs, but also from noise components, including detector dark counts, unsuppressed pump photons, uncorrelated SpRS in any silica components and multi-pair generation, collectively known as accidental counts. The accidental counts and genuine counts, which arise in the same pulse, cannot be deterministically separated; thus, one characterises the accidental counts by measuring subsequent pulse correlations with the original.

The silicon PhCWGs were fabricated from an SOI wafer, with a 220 nm-thick Si layer above a 2 µm layer of silica. The photonic crystal was patterned using electron beam lithography and reactive ion etching to create a triangular lattice of holes. The waveguides were made by introducing a row defect. The photonic crystal region was then undercut by etching away the silica substrate in that region, suspending the PhCWGs in air. The two rows of holes adjacent to the waveguide were laterally shifted to engineer the dispersion [154] so that the group index was ~ 30 across a 15 nm bandwidth, and centred at a wavelength of 1,559 nm. The effective nonlinearity $\gamma_{\text{eff}} = (n_g/n_0)\gamma$ was 4,000 W$^{-1}$m$^{-1}$, where the intrinsic material nonlinear coefficient, $\gamma$, the slowdown factor, $S$, and ratio of the group index, $n_g$, to the native refractive index of the material, $n_0$, is included. The PhCWG had a linear propagation loss of 50 dBcm$^{-1}$. The waveguides were 196 µm long, with inverse tapers and SU8 polymer cladding for improved coupling by mode matching to input-output lensed fibres. The first device measured had many separate PhCWGs with individual input and output coupling. The second sample included Y-split couplers, fabricated using silicon nanowires preceding the pairs of PhCWGs. The PhCWG regions were fabricated as close as possible to each other in order to avoid variations in the Si layer thickness across the surface of the wafer, thereby minimising any differences in the dispersion and group index.
4.2 Photon statistics measurements

This Section measures the photon statistics for two different photonic crystal device configurations. The first consists of two PhCWGs with the classical pump beam coupled separately to each waveguide on the chip, allowing for ease of balancing the coupled power. The second has an additional an on-chip Y-coupler to divide the input pump laser equally between two PhCWGs, demonstrating further integration.

4.2.1 Dual-input device

In the first demonstration, the two waveguides are labelled $A_1$ and $B_1$, and are imaged in Fig. 4.3. The rate of pairs from each source was measured before the electro-optic switch and compared with the multiplexed rate measured at the common output. The standard signal-to-noise metric for probabilistic photon-pair sources is the CAR. Figure 4.4 illustrates the CAR for a range of measured heralded single-photon rates for individual PhCWGs $A_1$ and $B_1$, shown as red squares and green circles, respectively. The characteristic curve

![Figure 4.3: a) Microscope image of input lensed fibres to two photonic crystal waveguides. b) SEM image of the photonic crystal waveguide. c) SEM image of the SU8 polymer mode converter.](image)
for the measured CAR of an individual source was limited by multi-pair noise. The absolute rate was then estimated by taking into account component losses and detector efficiencies. The two source were spatially multiplexed by adding an optical switch and electronics to route photons to a common output fibre. The measurements were made using the same characterisation setup as for the individual sources, and the result is plotted in Fig. 4.4 as blue triangles.

The data were fitted for the single sources using an analytical expression for the CAR ratio, including the measured values for detector dark counts and losses. By measuring the transmission of the opto-ceramic switch channels (85.1 and 79.4%), the maximum expected enhancement to the CAR was calculated as a function of the heralded single-photon rate (dashed line in Fig. 4.4, see Section 4.3 for more details).

4.2.2 Single-input device

The second experimental implementation demonstrates a primary building block for scalable spatial multiplexing. The two PhCWG devices are referred to
as ‘A2’ and ‘B2’, shown in the SEM image in Fig. 4.5a) and nominally identical dispersion plotted in Fig. 4.5b). The output heralded single-photon rate from each waveguide was measured and plotted in Fig. 4.6, shown as green triangles and red squares for A2 and B2, respectively. The multiplexing switch was then added and the multiplexed output measured for the same range of PhCWG output powers, shown in Fig. 4.6 as blue triangles. This study extracted an enhancement to the heralded single-photon rate of 63.1% (see Section 4.3 for more details), consistent with the result of the first measurement.

Figure 4.5: a) SEM of the single-input dual PhCWG device. b) Measurement of the matching dispersion curves for each device at 1,550nm.

Figure 4.6: Single-photon rates from source A2 (green circles), B2 (red squares) and the multiplexed (blue triangles) rate are plotted for a range of output powers.
## 4.2.3 Second order correlation measurements

To verify that the output of the multiplexed source was indeed in the single-photon regime, the second-order correlation function, $g^{(2)}(nT)$, was measured at discrete delays using a HBT setup [9,208]. Here, $C_{ABH}$ is a threefold coincidence across all detectors, $C_{AH}$ and $C_{BH}$ are the heralded clicks in detectors $A$ and $B$, $C_H$ is the number of heralding clicks, $n$ is an integer and $T$ is the period between pump-laser pulses. The following expression

$$g^{(2)}(nT) = \frac{C_{ABH}C_H}{C_{AH}C_{BH}},$$  \hspace{1cm} (4.1)
Figure 4.8: $g^{(2)}(0)$ experimental schematic: Two heralded single photon sources are spatial multiplexed to a common single output. The photons from this output, whose rate is enhanced compared to a single device, is then sent to the input of a 50:50 fibre coupler. This fibre coupler acts as the beam-splitter in a HBT experiment. The two outputs of the fibre coupler go to SPDs. Operation in the single-photon regime will result in very few coincidences being heralded across the output of the fibre coupler, corresponding to a $g^{(2)}(0) \ll 0.5$.

Figure 4.9: The $g^{(2)}(nT)$ correlation function is measured for the multiplexed output. A value of $g^{(2)}(0) = 0.19$, markedly $< 0.5$, confirmed that the source was operating in the single-photon regime. For other values of $n$, the uncorrelated Poissonian $g^{(2)}(nT)$ was close to 1, as shown by the dashed line. All errors were calculated from Poissonian statistics.
line (EDL) was added to adjust the fine balance in optical path length, with the requirement that the optical paths from each of the photon-pair sources to the coupler be balanced. At zero time delay between the detectors, a \( g^{(2)}(0) < 1 \) is expected for a non-classical light source, and \( g^{(2)}(0) < 0.5 \) for a source approaching single-photon operation.

The second order correlation function \( g^{(2)}(nT) \) was measured after multiplexing the photons from two PhCWGs using a HBT setup [17]. In this experiment, sources \( A2 \) and \( B2 \) both contribute to the measured photon statistics used to calculate \( g^{(2)}(nT) \), which is plotted in Fig. 4.9. The measured \( g^{(2)}(0) = 0.19 \pm 0.02 \) confirmed that the spatially multiplexed source was operating in the single-photon regime. To check this result, \( g^{(2)}(nT) \) was measured for \( n = \pm 1, 2 \), by delaying or advancing the electronic trigger pulses from the detector at one output of the 50/50 coupler by one or two pulse periods with respect to the other, and measuring the correlation. For neighbouring pulses, this was found to be unity within error, consistent with the Poissonian statistics expected when measuring distinguishable photons [131,177,209]. For these correlations between successive pulses, the photon statistics were expected to return to a Poissonian distribution with \( g^{(2)}(\pm nT) = 1 \). The measured \( g^{(2)}(+T) = 1.00 \pm 0.22, g^{(2)}(-T) = 0.91 \pm 0.21, g^{(2)}(+2T) = 1.14 \pm 0.26 \) and \( g^{(2)}(-2T) = 0.95 \pm 0.22 \) were consistent with this expectation. The data for all five peaks were taken for seven hours each, resulting in a total number of threefold counts for the \( g^{(2)}(0) \) case of 46, and 18 to 20 counts for all other cases. The data for the \( g^{(2)}(0) \) experiment are tabulated in Fig. 4.10.

<table>
<thead>
<tr>
<th>( g^{(2)}(nT) = (C_{ABH}/C_{A_{BH}}) )</th>
<th>( n )</th>
<th>-2</th>
<th>-1</th>
<th>0</th>
<th>+1</th>
<th>+2</th>
</tr>
</thead>
<tbody>
<tr>
<td>3-fold Coincidence</td>
<td>( C_{ABH} )</td>
<td>19</td>
<td>18</td>
<td>46</td>
<td>20</td>
<td>19</td>
</tr>
<tr>
<td>Coincidence Delayed</td>
<td>( C_{A_{BH}} )</td>
<td>19475</td>
<td>23096</td>
<td>280840</td>
<td>23213</td>
<td>23281</td>
</tr>
<tr>
<td>Coincidence Fixed</td>
<td>( C_{B_{BH}} )</td>
<td>544832</td>
<td>544832</td>
<td>544832</td>
<td>544832</td>
<td>544832</td>
</tr>
<tr>
<td>Heralds</td>
<td>( C_{H} )</td>
<td>5.78E+08</td>
<td>5.78E+08</td>
<td>5.78E+08</td>
<td>5.78E+08</td>
<td>5.78E+08</td>
</tr>
<tr>
<td>( g^{(2)}(nT) )</td>
<td>( g^{(2)}(nT) )</td>
<td>1.13759</td>
<td>0.908752</td>
<td>0.190989</td>
<td>1.004635</td>
<td>0.951616</td>
</tr>
<tr>
<td>Error ( g^{(2)}(nT) )</td>
<td>Err(( g^{(2)}(nT) ))</td>
<td>0.260981</td>
<td>0.214195</td>
<td>0.02816</td>
<td>0.224643</td>
<td>0.218316</td>
</tr>
</tbody>
</table>

Figure 4.10: Complete data for the correlation function measurement of the multiplexed stream of photons. All data sets were taken for seven hours in total.

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4.3 Analysing the multiplexing enhancement

In the previous sections, the experimental viability of multiplexing was demonstrated. This section first calculates the enhancement to the CAR by modelling a single source and the multiplexed source, including the switch efficiencies using a photon statistics method. The scalability of spatial multiplexing is then explored using an exact matrix method. For the experimental parameters used in the previous experiment, peak enhancement is achieved by multiplexing 128 sources together.

4.3.1 Photon statistics by approximate analytic method

This began by taking a single source, noting that true coincidences, $C^{(1)}$, per pulse generated through SFWM can be expressed as $C^{(1)} = \mu \eta_S \eta_I$, where $\mu$ is the number of pairs generated per pulse, and $\eta_S$ and $\eta_I$ are the overall collection efficiencies for the signal and idler photons, respectively. Here a lumped statistics approach is used and $\mu < 0.1$ ensured. The accidentals could then be calculated as the product of the singles rates, $N_{S,I} = \mu \eta_{S,I} + d_{S,I}$, such that $A^{(1)} = N_S N_I = (\mu \eta_S + d_S)(\mu \eta_I + d_I)$, where $d_S$ and $d_I$ are the dark counts in the signal and idler detectors, respectively. CAR$^{(1)}$ was calculated by taking the ratio of coincidences to accidentals such that

$$CAR^{(1)} = \frac{C^{(1)}}{(\mu \eta_S + d_S)(\mu \eta_I + d_I)}. \quad (4.2)$$

Then $\mu = C^{(1)}/\eta_S \eta_I$ could be substituted in the denominator to obtain an equation that was only dependent on $C^{(1)}$

$$CAR^{(1)} = \frac{C^{(1)}}{(\eta_S + d_S)(\eta_I + d_I)}. \quad (4.3)$$

For the two source multiplexed case, the coincidences could be expressed at the output of the switch as $C^{(2)} = \mu \eta_S \eta_I \eta_{L1} + (1 - \mu \eta_I - d_I) \mu \eta_S \eta_I \eta_{L2}$, whereas the accidentals were $A^{(2)} = A_1^{(2)} + A_2^{(2)} = (\mu \eta_I + d_I)(\mu \eta_S \eta_{L1} + d_S) + (1 - \mu \eta_I - d_I) \times (\mu \eta_I + d_I)(\mu \eta_S \eta_{L2} + d_S)$. Here, $\eta_{L1,L2}$ are the switch losses for the two sources, and it is assumed that the detector dark counts and idler
efficiencies are balanced for the two sources. Thus, \( \mu \) could be substituted as in the single-source example to obtain an expression for the enhanced CAR for two sources in terms of \( C^{(1)} \)

\[
CAR^{(2)} = \frac{[\eta L_1 + \eta L_2 (1 - C^{(1)}))] C^{(1)}}{(\frac{C^{(1)}}{\eta s} + d_I)(\frac{C^{(1)} \eta L_1}{\eta s} + d_S) + (1 - \frac{C^{(1)}}{\eta s} - d_I)(\frac{C^{(1)}}{\eta s} + d_I)(\frac{C^{(1)} \eta L_2}{\eta s} + d_S)}.
\] (4.4)

These equations were used to create a parametric plot of CAR versus \( C^{(1)} R \), where \( R \) is the repetition rate of the laser, as shown in Fig. 4.11. To calculate \( \eta_{S,I} \), the single-source measurement data, plotted as green circles and blue squares, were fitted using equation 4.3 with the measured dark counts per pulse \( d_S = d_I = 2 \times 10^{-6} \) and taking the average, shown as the blue dashed line, resulting in \( \eta_S = \eta_I = 0.001 \). Using these parameters in equation 4.4 and taking the experimentally measured switch losses of \( \eta_{L_1} = 0.79 \) and \( \eta_{L_2} = 0.85 \) the multiplexed curve was plotted as the red solid line, which is in agreement with the data shown as red triangles. The case for a lossless switch \( (\eta_{L_1} = \eta_{L_2} = 1) \) is shown as a black dotted line.

![Figure 4.11: Plot of the experimental data for the separately pumped PhCWG experiment showing CAR versus the heralded single-photon rate. The single source data are shown as blue squares and green circles, fitted with a dashed blue line. The multiplexed data are shown as red triangles, fitted with a solid red line. The theoretical lossless switch case is shown as a black dotted line. The error bars were derived from Poissonian statistics.](image)

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$\eta_{L2} = 1$) is plotted as a black dotted curve. The experimental enhancement can be extracted by taking the ratio of the multiplexed count output, $C^{(2)}$, and the single-source rate, $C^{(1)}$, calculated to be 62.4%.

To fit the data of coincidences versus power for the Y-split device, a similar method was used. Coincidences per second can be described by $C^{(1')}_1 = \mu \eta S \eta I R$, and the fact that $\mu = A(\gamma PL)^2$, where $P$ is the coupled power, $\gamma$ is the nonlinear parameter, $L$ is the length of the waveguide and $A$ is the SFWM efficiency. The single source data are shown as green circles and blue squares in Fig. 4.12. The fitting functions, shown as the blue dashed line, were used to find $RA(\gamma L)^2\eta S \eta I$. Taking the multiplexed output $C^{(2')}_1 = R\mu \eta S \eta I L_1 + R(1 - \mu)\mu \eta S \eta I L_2$ and inserting $RA(\gamma L)^2\eta S \eta I$ and the switch losses of $\eta_{L1} = 0.79$ and $\eta_{L2} = 0.85$, the multiplexed CAR curve was plotted as the red solid line, which was in agreement with the data shown as red triangles. The enhancement was the ratio of these curves, giving an enhancement of 63.1%. Finally, the case for a lossless switch with an enhancement of 100% was plotted as the black dotted curve.

![Figure 4.12: Plot of the experimental data for the Y-split device showing heralded single-photon rate versus output power. The single source data are shown as blue squares and green circles, fitted with a blue dashed line. The multiplexed data are shown as red triangles fitted with a red blue line. The theoretical lossless switch case is shown as a black dotted line.](image-url)
4.3.2 Photon statistics by matrix method

The model presented in this section is based on the method presented in reference [69], extended to include multiplexed sources. The experimentally observable properties in the model are the detector click probabilities, $\vec{P}_{\text{Clicks}}$. For example, $P_{HS}$ is the probability that the heralding arm has clicked and that the signal arm has not clicked. The detector click probability vector has the form

$$\vec{P}_{\text{Clicks}} = \begin{bmatrix} P_{HS} & P_{HS'} & P_{HS''} \end{bmatrix}^T.$$  \hspace{1cm} (4.5)

Initially none of the detectors has clicked and the probability vector is in the state

$$\vec{P}_{\text{Initial}} = \begin{bmatrix} 1 & 0 & 0 \end{bmatrix}^T.$$  \hspace{1cm} (4.6)

The probability that any given detector will click, per single pump laser pulse, can be calculated by acting on $\vec{P}_{\text{Initial}}$ with a matrix corresponding to the probability of a photon pair click, taking into account losses. This matrix has the form

$$M_{\text{Loss}} = \begin{bmatrix} (1 - \eta_H)(1 - \eta_S) & 0 & 0 & 0 \\ \eta_H(1 - \eta_S) & (1 - \eta_S) & 0 & 0 \\ (1 - \eta_H)\eta_S & 0 & (1 - \eta_H) & 0 \\ \eta_H\eta_S & \eta_S & \eta_H & 1 \end{bmatrix},$$  \hspace{1cm} (4.7)

where $\eta_H, \eta_S$ are the collection losses in the herald and signal arms, respectively. $\eta_{H,S}$ and $(1 - \eta_{H,S})$ are the probabilities that the herald/signal detector has clicked or not clicked, respectively. Dark counts can also cause a detector to click, and are also included with the form

$$M_{\text{DC}} = \begin{bmatrix} (1 - D_H)(1 - D_S) & 0 & 0 & 0 \\ D_H(1 - D_S) & (1 - D_S) & 0 & 0 \\ (1 - D_H)D_S & 0 & (1 - D_{\text{MUX}}) & 0 \\ D_HD_S & D_S & D_{\text{MUX}} & 1 \end{bmatrix},$$  \hspace{1cm} (4.8)

where $D_{H,S}$ are the probability of detector dark clicks per pump pulse. The
final detector states, $\vec{P}_{\text{Final}}$, were calculated using the following expression [69]

$$\vec{P}_{\text{Final}} = \sum_{i=0}^{N} p_{\text{thermal}}^i \cdot M_{\text{DC}} \cdot (M_{\text{Loss}})^i \cdot \vec{P}_{\text{Initial}}, \quad (4.9)$$

where $P_{\text{thermal}}^i$ is the probability that $i$ photon pairs were generated for a given $\mu$, average pairs generated per pump pulse. A thermal distribution of the number of photon pairs generation is described by the following expression

$$p_{\text{thermal}}^i = \frac{\mu^i}{(1 + \mu)^{i+1}}. \quad (4.10)$$

A thermal distribution was assumed for this calculation because the nonlinear photon sources investigated in this chapter had a photon coherence length, $\tau_{\text{Photon}}$, close to the pump pulse period, $T_{\text{Pump}}$. Details on when and why to use a thermal distribution or Poissonian distribution for photon statistics calculations were presented in Section 2.2.1.

**Active spatial multiplexing of high loss sources**

Now, the above method is extended to include spatial multiplexing of lossy sources. This method holds for devices where the multiplexed source never reaches the performance of a single ideal, lossless nonlinear photon source. This is valid when dealing with on-chip generation where coupling and propagation losses are intrinsically high. The specific multiplexing scheme is detailed in Fig. 4.13. For these calculations, the number of sources can only increase in powers of two—that is, one source, two sources, four sources and so forth. Each jump in the number of sources corresponds to an additional layer of multiplex switches.

The herald photon from each source is immediately detected by one of $N$ detectors. The sources are ordered in priority from one to $N$, with the lowest loss source assumed to be in the highest priority location. Concretely, if $\eta_H$ is the probability that detector 1 is triggered and $(1 - \eta_H)$ that it does not click, then the probability that detector 2 clicks is $\eta_H(1 - \eta_H)$. The total probability
of a herald detector making a photon detection is then

\[ \eta_{H_{\text{mux}}} = \sum_{i=1}^{N} \eta_H (1 - \mu \eta_H)^{N-1}, \]  \hspace{1cm} (4.11)

where N is the total number of multiplexed sources. It is important to note that once a herald detection has been made, the remaining herald detections are ignored. Each additional heralding arm added for multiplexing has a probability of contributing a heralding click due to dark counts. The dark count contribution for the heralding arm then must include all N detectors,

\[ D_{\text{mux}} = \sum_{i=1}^{N} D_H (1 - D_H)^{N-1}, \]  \hspace{1cm} (4.12)

Once again, if a detector clicks due to a dark count, the remaining heralding clicks are ignored. The loss in the signal arm now must take into account the
loss of each switching layer, with the total signal loss equal to

\[ \eta_{\text{Smux}} = \eta_S (\eta_{\text{switch}})^{\log_2(N)}. \]  \hspace{1cm} (4.13)

This expression assumes a non-zero switch loss of \( \eta_{\text{Switch}} \) combined with the intrinsic collection loss in each of the signal arms, \( \eta_S \). Now to calculate the photon statistics for a multiplexed source, it is necessary to apply the new probability matrices for the multiplexed losses,

\[ M_{\text{Spatial}} = \begin{pmatrix}
(1 - \eta_{\text{Hmux}})(1 - \eta_{\text{Smux}}) & 0 & 0 & 0 \\
\eta_{\text{Hmux}}(1 - \eta_{\text{Smux}}) & (1 - \eta_{\text{Smux}}) & 0 & 0 \\
(1 - \eta_{\text{Hmux}})\eta_{\text{Smux}} & 0 & (1 - \eta_{\text{Hmux}}) & 0 \\
\eta_{\text{Hmux}}\eta_{\text{Smux}} & \eta_{\text{Smux}} & \eta_{\text{Hmux}} & 1
\end{pmatrix}, \]  \hspace{1cm} (4.14)

and the multiplexed dark counts,

\[ M_{\text{MUXDC}} = \begin{pmatrix}
(1 - D_{\text{mux}})(1 - D_S) & 0 & 0 & 0 \\
D_{\text{mux}}(1 - D_S) & (1 - D_S) & 0 & 0 \\
(1 - D_{\text{mux}})D_S & 0 & (1 - D_{\text{mux}}) & 0 \\
D_{\text{mux}}D_S & D_S & D_{\text{mux}} & 1
\end{pmatrix}. \]  \hspace{1cm} (4.15)

Equation 4.16 subsequently gives the probabilities for an actively multiplexed source,

\[ \tilde{P}_{\text{Final}} = \sum_{i=0}^{N} p^{i}_{\text{thermal}} M_{\text{MUXDC}} (M_{\text{Spatial}})^i \tilde{P}_{\text{Initial}}. \]  \hspace{1cm} (4.16)

Useful values can be extracted from \( \tilde{P}_{\text{Final}} \) for the coincidence and singles rates. The single-photon count rates for the herald and signal are \( N_S = P_{\text{HS'}} + P_{\text{H'S'}} \) and \( N_H = P_{\text{H'S}} + P_{\text{H'S'}} \), respectively. The total coincidence rate is given by \( C_{\text{total}} = P_{\text{H'S'}} \) and the accidental coincidence rate \( A = N_H \times N_S \). The true coincidence rate, attributed to the nonlinear pair generation is then \( C_{\text{true}} = P_{\text{H'S'}} - A \). Finally, the CAR is as defined in Section 2.2.1, with \( \text{CAR} = (C_{\text{total}} - A)/A \).
Figure 4.14: a) CAR plots for a collection efficiency of -23 dB and 1 dB switch loss for one, two, and four sources, as well as the CAR plot for eight multiplexed sources with no switch loss. b) Maximum CAR for increasing number of multiplex sources, for 0, 0.5, 1 and 1.5 dB switch loss, respectively.

Practical limits to active spatial multiplexing

To explore the practical limits of spatial multiplexing, the methodology introduced in the previous section was used with realistic experimental parameters, setting the collection loss in both arms to -23 dB, $\eta_S = \eta_H = 0.005$. The dark counts were set to 500 Hz with an operating repetition rate of 50 MHz, setting $D_H = D_S = 10^{-5}$. CAR versus $C_{true}$ is plotted in Fig. 4.14. Multiplexing additional sources improves the heralded single-photon rate (coincidences between idler and signal photons), but cannot increase the maximum CAR. An upper limit on the CAR of ~125 was set by the collection and detection efficiency—the dashed line in Fig. 4.14a).

In addition, the multiplexing switch loss was set to the conservative value of -1 dB, $\eta_{\text{switch}} \sim 0.79$. These were the approximate parameters used in the experiment of Section 4.2.2. Figure 4.14a) aligns well with Fig. 4.11 from Section 4.3, plotted using an analytical method. For a fixed CAR below the maximum, the photon heralded single-photon rate was enhanced with every additional photon pair source added. Even for a lossless switch, the additional dark counts reduced the maximum CAR, shown in Fig. 4.14b) as the gold dots. Figure 4.14b) shows how the maximum achievable CAR dropped with additional sources, for four different switch losses.

Figure 4.15 plots the maximum heralded single-photon rate at the minimum useful CAR threshold of 20. Figure 4.15a) displays the experimental parameter
of collection loss $\eta_S = \eta_H = -23dB$ and switch loss $\eta_{\text{switch}} = -1dB$, again close to the experiment of Section 4.2.2. The optimum number of multiplexed sources for this experimental regime was 128, with negligible improvement gained by going to 256 (requiring double the physical resources). The 128 multiplexed source would be 20 times brighter than an individual source, when operating at a CAR of 20. Reducing the switch loss to $\eta_{\text{switch}} = -0.5dB$, increased the heralded single-photon rate by $\Delta n_{\text{switch}}$, as shown in Fig. 4.15b). The optimum number of sources to multiplex then shifted to the larger number of 256.

This is now, contrasted to the statistics for a multiplexed bulk source with lower collection loss of $\eta_S = \eta_H = -10dB$, such as a lithium niobate, $\chi^{(2)}$ nonlinear waveguide [198]. In Fig. 4.15c), the switch loss is $\eta_{\text{switch}} = -1dB$

![Figure 4.15: The heralded single-photon rate at a CAR of 20 for increasing number of sources. The x-axis is plotted in a log_2 scale. a) For a collection loss of $-23dB$ and switch loss of $-1dB$, the peak heralded single-photon rate occurs for 128 multiplexed sources. b) For a collection loss of $-23dB$ and switch loss of $-0.5dB$, the peak heralded single-photon rate occurs for 256 multiplexed sources. c) For a collection loss of $-10dB$ and switch loss of $-1dB$, the peak heralded single-photon rate occurs for 16 multiplexed sources. d) For a collection loss of $-10dB$ and switch loss of $-0.5dB$, the peak heralded single-photon rate occurs for 16 multiplexed sources.](image-url)
and the optimum heralded single-photon rate operating at a CAR of 20 is reached after multiplexing 16 sources. The heralded single-photon rate of this multiplexing source is $\sim 10^3 \times$ brighter than the previous source used to calculate the plot in Fig. 4.15a). Reducing the loss of the switch to $\eta_{\text{switch}} = -0.5 \text{dB}$, Fig. 4.15d), provides a modest improvement to the heralded single-photon rate. The calculations presented in this section show the potential for effectively scaling up the multiplexing scheme presented earlier in this chapter, to a maximum of 128 sources. It is also indicates that improving the quality of each individual source is vital to reduce the number of sources required to reach the optimum heralded single-photon rate using multiplexing.

### 4.3.3 Discussion

The efficiency of the heralding detectors used here was limited to 10% when operating at 1,550 nm in order to limit dark counts to reasonable levels. In a multiplexed source, any improvement to the detector or heralding efficiency carries over favourably to the single-photon output rate. For example, doubling the detection efficiency for a source with N multiplexed waveguides would result in a $2^N$ increase in the heralded single-photon rate.

The linear losses associated with the waveguide and filtering components will likely be reduced in the future with advances in fabrication technology, thereby further increasing the single-photon rate. Monolithic MZI arrays, realised using low-loss thermal switches in silicon, may prove to be a preferable technology when scaling to large numbers of switches, because the input and output insertion loss does not change for one switch or a whole array. Although the current demonstration of spatial multiplexing implemented here is not monolithically integrated onto a single photonic chip, it is all-fibre coupled, making it a stable and robust platform.

A major hurdle in moving to a monolithic device is the development of low-loss optical delay lines (ODLs) that are compatible with low-noise photon pair sources, preferably in a single material platform, which is currently not the case. Recent integrated sources show superior output photon statistics, to those presented here, including heralding efficiencies [121]. However, this study stresses that this multiplexing scheme is applicable to any of these sources, and
will provide comparable improvement for any source.

The results presented in this chapter show that integrated spatial multiplexing can be implemented to efficiently improve the statistics of a nonlinear photon source. An enhancement of 62.4 and 63.1% to the probability of generating a single-photon was achieved, thereby breaking the intrinsic limit of a single source by decoupling the single and multi-pair noise probabilities. In addition, the building-block demonstration has established the feasibility of scalable multiplexed sources. Multiplexing will continually benefit from technological development in the areas of ultrahigh efficiency detectors, low-loss integrated optics and precision fabrication, and will eventually lead to impressive multi-qubit (> 10) demonstrations.
Chapter 5

Conclusions, perspective and future directions

This chapter presents a summary and concluding remarks on the original work presented in this thesis, in addition to discussing a perspective on the field and its future direction.

‘Mr. Data, lay in a course for the 24th century. I suspect our future is there waiting for us.’

– Captain Jean-Luc Picard
5.1 Summary

Photonic technology has provided a mature experimental toolbox for QIP using single-photons. However, as discussed in Chapter 2, many fundamental challenges associated with efficient single-photon generation remain. One solution is to use intrinsic material nonlinearity for single-photon generation. This thesis set out to answer two key questions regarding nonlinear single-photon generation in photonic waveguides:

1. Is chalcogenide glass a viable platform for on-chip nonlinear single-photon generation?

2. Can multiple heralded single-photon sources be combined using active spatial multiplexing to improve the overall photon statistics?

The first question was posed as a result of the original chalcogenide photon pair generation experiment in 2011 [86]. In reference [86] the coincidence to accident ratio (CAR, a measure of the photon-pair signal-to-noise ratio) was measured to be 0.08–far below the useful minimum of 10. It was suggested, but not confirmed, that this was due to uncorrelated Raman photons generated in the waveguide. Chapter 3 began with the initial intuition that reducing the phonon population by cooling would eliminate Raman noise. Cooling a chalcogenide fibre to 77 K in liquid nitrogen reduced near pump Raman noise by a factor of $\sim 3$, a modest improvement. Moving from a CW to pulsed pump further improved the CAR by an order of magnitude. These two points together increased the CAR to four, still below the usable limit.

Measuring the photon pair correlation is an indirect method of observing the effect cooling has on the Raman scattering. A direct measurement method was needed, which motivated the design of a novel single-photon counting Raman spectrometer. This was used to directly measure the Raman spectrum in chalcogenide and other photonic waveguide materials, both at room temperature and at 77 K [182]. The shape of the Raman spectra could now accurately be mapped and the effects of cooling directly monitored. The spectra confirmed both the modest effects of cooling and the existence of an intrinsic minimum at $7.4 \text{ THz}$ from the pump, dubbed the low-Raman window for chalcogenide.
To target the ‘low-Raman window’ and have the signal-photon generated at \( \sim 1.55 \mu m \), a dispersion-engineered waveguide was required. The dispersion of the TM mode of a ridge waveguide would allow for the desired pair generation bandwidth. However, ridge waveguides are designed for TE polarisation operation, with high losses in the TM polarisation. By using a short device, the excess loss accrued from TM operating was minimised, allowing efficient pair generation. The CAR improved to 16.8 in the low-Raman window, a 25 times improvement on the near-detuned case. It was conclusively shown that Raman noise is detrimental to single-photon generation in the chalcogenide platform, across a broad bandwidth, with the exception of the low-Raman window.

The first multiplexing demonstration used bulk nonlinear crystals [134] and faced problems of stability that limited the scalability of the approach. These issues can be overcome by moving the multiplexing to an integrated platform. In Chapter 4, the first active spatial multiplexing of integrated heralded single-photon sources was demonstrated. This was realised by overcoming several technological challenges. First, two slow-light photonic crystal waveguides with nominally identical dispersion and a Y-split input were successfully fabricated. Fitting the whole structure in a single electron beam write field ensured that the exposure dose and alignment was the same for both devices. In addition, any variations in the thickness of the silicon wafer were avoided because these occur on a length scale larger than the 10 \( \mu m \) waveguide separation.

Next, the combined loss of the ODLs, fibre connectors and optical switches needed to be significantly less than 50\% to see a compelling net gain. The use of ultra-low-loss PLZT electro-optic modulators and careful optimisation of the setup made this possible. Finally, broadband electronic analogue logic was needed to trigger the PLZT modulators at the 100 ns time scale required. Combining these culminated in a 60\% multiplexing enhancement to the useful heralded single-photon rate. Theoretical analysis showed that the enhanced photon statistics could not be reached by a single source. Further, it was shown that spatial multiplexing scales favourably up to 128 similar sources. These experiments conclusively showed that combining multiple heralded single-photon sources improves the overall photon statistics. This work was the first step towards a scaled multiplexed scheme containing many sources, whose enhanced photon statistics will unlock a new regime of QIP protocols.
5.2 Chalcogenide in quantum photonic networks

At first glance, chalcogenide appears to have many of the properties desired for photonic nonlinear single-photon generation. Chalcogenide can be fabricated into planar waveguides, and, due to having a high refractive index, allows tight waveguide bend radii and thus complex circuits. Chalcogenide is highly nonlinear, thereby enabling photon generation in short device lengths. However, like all amorphous materials, the Raman noise in chalcogenide is generated over a broad bandwidth, overlapping the useful photon generation channels, with the exception of an intrinsic minimum in the Raman spectrum that lies at 7.4 THz detuned from the pump. Operating in the low-Raman window requires a herald and signal photon separation of 14.8 THz to maintain energy conservation. This means that one of the photons must lie far outside the telecommunication range, increasing the complexity of the source.

While chalcogenide is not a platform suited for single-photon generation, there are other applications in photonic quantum technology where chalcogenide may yet take the lead. One specific example is wavelength agility in single-photon based quantum communication networks [187]. To satisfy the flexibility and bandwidth demand of modern telecommunication, networks are agile and reconfigurable. Reconfigurable optical add-drop multiplexers [174] and tuneable transmission lasers provide network reconfigurability and bandwidth flexibility, respectively. To implement the same functionality in a quantum network requires the ability to convert single photons from one wavelength to another, while retaining the photon’s quantum information. This can also be used to adapt the senders and receivers requirements that, until a clear winning technology emerges, may not operate at the same wavelength [210,211].

Frequency conversion in the single photon regime, using Bragg-scattering four-wave-mixing [89], was demonstrated in 2012 using silica fibre with 99% conversion efficiency [178]. Two challenges facing this experiment were the long length of fibre required (∼800 m) and the effects of Raman depletion. A short centimetre long chalcogenide waveguide could be used to achieve the same level of nonlinearity required in the 2012 experiment, while operating in the low-Raman window would eliminate both the problems of spontaneous Raman noise and stimulated Raman depletion. To avoid both of these problems in
silica would require operating at the impractical detuning of > 15 THz, beyond the Raman band in silica. The inset in Fig. 5.1 shows the measured SpRS, using the technique from Section 3.2. This shows that no intrinsic minimum exists for silica. Chalcogenide waveguides can also be dispersion engineered, with the zero dispersion wavelength, $\lambda_0$, a key parameter for frequency conversion [212]. Silicon could not be substituted for chalcogenide because TPA would limit the pump power before 100% conversion could be achieved, and the free carriers generated from TPA add significantly to the nonlinear losses. Another quantum protocol requiring strong nonlinearity is coherent photon conversion [186], where chalcogenide has been flagged as a well suited platform. In summary, chalcogenide integrated circuits are set to play a major role in frequency conversion, a vital tool in future agile quantum networks.

**Figure 5.1:** Illustration of the single-photon frequency conversion spectrum. The Raman noise generated by the strong pumps can be avoided by operating them at a detuning of $\sim$7.4 THz from the single-photon channels.
5.3 Multiplexing of heralded single-photon sources

This thesis has successfully demonstrated that spatial multiplexing enhances the photon statistics of lossy nonlinear, heralded single-photon sources [138]. This was enabled by key experimental developments, including near-identical photonic crystal waveguide fabrication, high bandwidth analogue electronic logic and ultra-low-loss PLZT optical switches.

In the future, these switches could also be implemented in silicon, with thermal phase modulators now having the speed and low-loss properties required [213]. The key challenges facing further scaling of an all-silicon multiplexed source are physical resources and low-loss delay lines. The first can be overcome by improving the quality of the photon sources. Section 4.3.2 indicated that, by reducing the photon collection/detection loss to 10%, the optimum number of multiplexed sources reduces to only 16.

The latter problem of delay is more fundamental and, for the near future, will halt the complete integration of a complex multiplexed source. The delay line must be long enough to give time for photon detection, electronic signal generation, radio frequency logic decision and switch-state change, while having low enough propagation loss to make multiplexing viable. A delay of 10 ns corresponds to an almost one metre long silicon waveguide. Currently, the propagation loss for a silicon waveguide is around $-1 \text{dB/cm}$ or $20\%/\text{cm}$. Excess loss renders all-silicon multiplexing untenable. However, recently, ultra-low-loss delay lines were demonstrated using exotic silicon waveguides [206].

![Spatial multiplexing scheme](image)

*Figure 5.2: Spatial multiplexing scheme. Ideally, to avoid losses coupling loss, all of the components depicted need to be fabricated in a monolithic silicon chip.*
5.3.1 Alternate spatial multiplexing schemes

In 2012, a scheme for asymmetric spatial multiplexing (ASMUX) was proposed [140], claiming to offer improved scaling over the original symmetric spatial multiplexing (SMUX) scheme [132]. This proposal attempted to use the intrinsic asymmetry of SMUX as a resource, illustrated in Fig. 5.3a). SMUX, as formulated in Section 4.3, has an inbuilt hierarchy, where the heralding detection from the lower ranked sources is only considered if the higher ranked sources have not triggered. The analysis in [140] shows how ASMUX asymptotically approaches a limit based on the switching and collection efficiencies. This, in contrast with SMUX, where an optimum number of sources is reached, followed by a reduction in performance for any additional sources.

For the regime of a large number of multiplexed sources, ASMUX appears to out-perform SMUX. However, this analysis does not include dark counts or multi-pair terms beyond the first order. These terms become very significant in the regime where the ASMUX method out-performs the SMUX method, thereby weakening the claim.

It should also be noted that an exciting multiplexing scheme based on quantum memories was recently proposed [214]. This scheme has the potential to enhance the photon statistics of heralded single photon sources by orders of magnitude, hinged on the efficiency, $\eta$, and time-bandwidth product, $B$, of the quantum memories used.

Figure 5.3: a) Schematic representation of ASMUX [140]. b) Cartoon illustration of how temporal multiplexing collapsed multiple time-bins into one, thereby increasing the probability of a photon detection [142].
5.3.2 Temporal multiplexing

Active temporal multiplexing was first proposed in 2011 by Jake Mower and Dirk Englund at the Massachusetts Institute of Technology [142]. Temporal multiplexing deterministically combined photons generated in different time-bins, into a target time-bin, shown schematically in Fig. 5.3b). multiplexing scales more favourably than spatial in regard to switch resources, with $\log_2(N + 1)$ switches required for an $N \times$ scale up instead of $(N - 1)$ switches, respectively. Temporal multiplexing only requires one nonlinear element to generate the photon pairs, automatically achieving spectral indistinguishability between photons. Note that dispersion effects due to path differences must be carefully managed to maintain this indistinguishability.

While temporal multiplexing is very promising, the technical challenges with respect to experimental implementation are significantly greater than for spatial multiplexing. For example, when spatial multiplexing 16 sources the electronic logic that triggers the electro-optic switches must have a timing resolution of approximately the laser repetition rate of $\sim 10$ ns. The electronic pulses sent from the detector need only be slightly shorter than this for error free operation. In contrast, for temporal multiplexing, this 10 ns is broken into $N$ time-bins. For 16 sources, each bin is less than 1 ns. For error-free operation of analogue electronic logic, the bandwidth of the electronic logic needs to be around 10’s of Gigahertz to support sub-nanosecond logic. The upper limit on electronic bandwidth, practically speaking, is 64 GHz. This allows pulses with rise times below 100 ps. This does not take into account the fact that even for single element radio frequency transistors or diodes, the components of even the most lean logic gate, have a propagation time of a few nanoseconds. These considerations limit the number of time-bins that can be used in practice.

Given that the optimum number of time-bins to be used is the same as the number of nonlinear elements to multiplex for the spatial scheme, then the most efficient scheme would employ a hybrid approach. In this approach, temporal multiplexing would be used up to the practical electronic limit, then spatial would be used for the remaining sources. This would require the least amount of physical resources, while still reaching the optimum multiplexing point.
5.4 Silicon quantum photonic technology

One of key goals of the field of quantum photonics, is to create a fully integrated quantum optical processor. Integration provides the stability and scalability required for complex circuits and, if fabricated in a highly nonlinear material such as silicon, single-photons can be generated on-chip. Silicon photonics has long been an exciting area of classical photonics research, driven by the powerful potential to piggy-back on the already mature CMOS fabrication technology used for electronic integrated circuits. Access to large commercial CMOS foundries for photonics research has become widely available via multiple user programs. This has made ubiquitous the use of silicon photonics in quantum labs, and accelerated key component development. As of the conclusion of this thesis (late 2014), silicon photonics is poised to have a huge influence on the status-quo for quantum technology over the next five years.

The KLM scheme for LOQC requires low-loss linear optical components, single-photon sources and efficient SPDs. Heralded single-photon sources have been demonstrated with great maturity on-chip. Compact resonant [215] and non-resonant broadband [86] heralded photon sources have been demonstrated in silicon, as well as degenerate photon pair generation [216, 217]. These previous demonstrations purified the frequency spectrum with high extinction using bulk filters located off the chip. Recently, high quality on chip filtering was demonstrated [218], thereby completing the tool-box required to complete monolithic on-chip heralded single-photon generation.

Silicon waveguides losses have been steadily decreasing for the last decade and are now low enough for complex linear circuits. Active components are usually much less efficient, with high-speed modulators typically having only 10 to 20% transmission [219]. In 2014, low-loss phase shifters have been demonstrated in silicon [213]. In a MZI configuration, these can be used to make four-port switches and variable power couplers. These low-loss tuneable MZIs can be used to make arbitrary and reconfigurable silicon quantum circuits [220].

High-efficiency SPDs that operate by detecting photons evanescently inside a silicon waveguide have been successfully demonstrated for a number of years. However, the main challenge has been the low yield from the fabrication process. A scaled quantum processor will require a large number of detectors
to be successfully fabricated on a single chip. In 2014, a near-perfect yield method for fabricating high-efficiency photon detectors on silicon was demonstrated [221]. Note that these detectors operate at cryogenic temperatures (∼1 K).

All of the components shown in Fig. 5.4, when brought together, complete the toolbox for a silicon quantum processor. Within the next few years, researchers will have the pleasure of being involved in many ground-breaking demonstrations of quantum technology, based on silicon photonics.

Figure 5.4: Components of future all-silicon photonic quantum processor. a) Heralded single-photon source, based on a photon crystal [138]. b) Programmable quantum processor, made from cascading MZI’s [220]. c) Near-perfect yield SPDs, fabricated on top of a silicon waveguide [221].
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