Breaking the restrictions of nearest neighbour interactions in mass-manufacturable silicon photonics: Applications in quantum information systems

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Declaration

I certify that except where due acknowledgement has been made, the work is that of the author alone; the work has not been submitted previously, in whole or in part, to qualify for any other academic award; the content of the thesis is the result of work which has been carried out since the official commencement date of the approved research program; any editorial work, paid or unpaid, carried out by a third party is acknowledged; and, ethics procedures and guidelines have been followed.

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Abstract

This doctoral thesis presents a scalable mass-manufacturable waveguide platform that can make use of an optical bus to provide an on-chip long-range communication channel. This optical bus is combined with an adiabatic transfer technique to provide robust and highly coherent transfer between waveguides. The potential for this robust transfer is explored in the context of sensitive quantum information processing and the designs for several important quantum gates are introduced. The addition of the optical bus allows multi-dimensional coupling that is not possible within current planar waveguide platforms.

Shallow ridge waveguides are fabricated using a standard silicon photonics platform. The propagation loss due to the lateral leakage was characterised over a range of waveguide widths and operating wavelengths. Losses as low as 0.087 dB/mm can be achieved by considering the lateral leakage effect during waveguide design.

A new approach to long-range coupling is described that combines the lateral leakage effect observed in shallow ridge waveguides with an adiabatic transfer technique. This platform enables the transport of light between two waveguides over surprising distances using radiation within the silicon slab as an optical bus. Due to the nature of the adiabatic protocol, the bus is minimally excited and the transport is highly robust. The ability of the bus is further demonstrated by introducing an intermediate waveguide between the pair that can be isolated and completely bypassed from the interaction.

The adiabatic optical bus is then extended to quantum information applications that can perform operations on photonic qubits. Several important gate designs are described that can produce a Hadamard, 50:50 and 1/3:2/3 beam splitter and a non-deterministic controlled NOT gate, with calculations showing one and two-photon gate operation. This is the first adiabatic gate demonstration that required a quantum description for photons.

In summary, this thesis introduces an optical bus that is suitable for fabrication in mass-manufacturable silicon integrated photonic waveguides. The application of an adiabatic protocol ensures the robust transfer of information using this bus. The high fidelity transport can be exploited for the development of integrated photonic quantum gates.
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Chapter 1

Introduction

There are many important scientific calculations that are either infeasible, or simply not possible using classical computational schemes. Quantum computational approaches have been theoretically shown to be useful across a range of such problems. These problems broadly fall into three classes: Hidden subgroup, of which factorisation is the main example [1], with direct applications in quantum cryptography [2, 3, 4], search problems (including database search [5]) and quantum simulation [6, 7].

Driven by these non-classical applications, the scientific community is actively searching for solutions that can enable practical quantum computing. The available platforms for quantum experiments can either be matter based, such as manipulating trapped ions [8], molecular vibrational states [9] and Bose-Einstein condensates [10], or they can be optically based [11]. Optically based platforms are proving attractive in terms of the high coherence and low loss properties of photons compared to matter based systems [12].

The traditional approach to quantum processing is to prepare a large scaled experiment that can carefully control and manipulate the state of qubits using free space optics, filters and beamsplitters [13, 14, 15]. A demonstration of Shor’s powerful factoring algorithm is presented in a free-space quantum optics experiment [14].

An ideal quantum computer should be scalable to the levels of complexity and capabilities of a modern classical computer, however operating on quantum information [12]. The requirement for large scale qubit applications suggests integrated photonics as the correct platform.

Integrated photonic technology has enabled these quantum experiments to be shrunk from large optical benches full of components onto a single chip [16, 17]. Integrated photonics can be realised through use of the mass-manufacture infrastructure of the integrated electronics industry, provided that it adheres to certain constraints. Integrated photonic devices have demonstrated a two photon
controlled-NOT gate operation [16, 18] and Shor’s factoring [19]. This is based on linear optics quantum computing, which is a non-deterministic method to achieve scalable but probabilistic two photon gates [11]. This thesis does not address scalable quantum computing, it is focussed on one set of elements that may be a part of a scalable quantum computer.

A key issue with quantum systems is the fragility of quantum states [12]. This is especially a concern when scaling to large numbers of gate operations. Adiabatic transfer methods are useful in conserving coherence and allow the robust transport of quantum information [20]. However, the operation of adiabatic gates differ from their non-adiabatic counterparts. The most basic adiabatic gate requires a three dimensional parameter space to provide a non-trivial geometric or Berry phase rotation [21, 22]. Unanyan, Shore and Bergmann introduced their approach to robust adiabatic geometrical gates using a four-state system with one shared excited state [21, 23, 24].

Increasing the dimension of the qubit to a three state superposition is called a qutrit. In addition to being a requirement for an adiabatic gate, access to higher order interactions can provide redundancies for quantum error correction schemes [25, 26] and also help in reducing the number of qubits or gates required for non-adiabatic operations [27]. For example, a 5 control gate Toffoli operation requires 64 two-photon gates \(2^{(n+1)}\) [28]. This can be reduced to 9 gates \((2n - 1)\) by accessing one higher order superposition ‘qudit’ of order \((n + 1)\) [29]. A qutrit gate has been demonstrated using photonic waveguides, but relies on the use of a third spatial dimension [30, 31] or exploitation of polarisation as an additional dimension [27] due to the restrictions imposed by nearest neighbour interactions. Further extensions in dimensionality are hard to conceive, particularly while remaining within the constraints required for planar mass fabrication. Quantum theory has suggested other uses and applications for multi-dimensional channel mixing [32, 33], which can make use of scalable \(N \times N\) interacting system [34].

This thesis aims to discover a scalable waveguide platform suitable for mass-manufacture. This platform must maintain the coherence of entangled photon states across many gate operations to be suitable for quantum information processing. Breaking free from nearest neighbour limitations would allow arbitrary waveguides on a common chip to interact.

1.1 Vision

Integrated photonic waveguides have emerged as a strong candidate to be the platform of choice for the next generation quantum computer. Recent high profile quantum experiments have shown calculations that are not feasible in a classical sense such as boson sampling [35, 36].
Figure 1.1: (a) Two waveguides can couple with a strength that is dependent on their proximity, (b) introducing a third waveguide makes it no longer possible to mutually couple the outer waveguides, (c) this restriction can be overcome by fabricating in a third dimension, which is not a CMOS compatible process, (d) the extra dimension only buys an extra level of coupling, with higher dimensional couplers not available due to nearest neighbour interactions.

Still, these ground-breaking results are based on $2 \times 2$ directional couplers. There is an interest in increasing this to a more arbitrary $N \times N$ channel system [32]. This is problematic for planar integrated photonic circuits as they rely on evanescent coupling to allow a photon to tunnel, or interact with another waveguide (Fig. 1.1(a)). This incurs a limitation on the waveguide arrangements, only allowing maximum of two channels to mix when planar, and these channels must be next to each other (Fig. 1.1(b)). To transport a photon across a large network of waveguides, it is necessary to implement sequences of SWAP gate operations which increase the complexity and fragility of the system. It is possible to introduce mutual evanescent coupling between a third waveguide by introducing a third dimension (Fig. 1.1(c)). However, nearest neighbour interactions limit more than three waveguides talking in a mutual arrangement (Fig. 1.1(d)).

To break away from these topological constraints, while remaining in a mass-fabricable planar topology, a long range coupling mechanism is required, one that can provide coupling between an arbitrarily scalable array (Fig. 1.2(a)). This should include the possibility of completely bypassing any intermediate waveguides that may be providing independent functionality elsewhere, while remaining
Figure 1.2: (a) Hypothetical waveguide platform that can provide the selective transport of light between a set of waveguides using a photonic bus, (b) the extension of this concept to an $N \times N$ array of waveguides that could operate as an arbitrary quantum unitary transformation on a photonic $N$-state qudit.

within the same integrated circuit. Providing stringent control over the coupling and taking care to preserve coherence could then make this a suitable platform for manipulating and transporting quantum information photons, such as an arbitrary $N \times N$ beam splitter which can provide a multi-dimensional unitary operation [11, 37] (Fig. 1.2(b)).

The primary objective of this research is to discover a way to achieve long-range selective communication between waveguides that is no longer solely restricted to nearest neighbour interactions.

### 1.2 Adiabatic passage as a transfer protocol

The previously defined specifications require a mass-manufacturable waveguide platform that can allow many waveguides to communicate through a common state as an optical bus. Such an optical bus will only be suitable for quantum information experiments if transport can be achieved in a fault tolerant manner.

There is a long history in the literature of techniques proposed to provide error free population transfer. It will be beneficial to review the evolution of this field
and one of these techniques is STImulated Raman Adiabatic Passage (STIRAP) [38].

### 1.2.1 STImulated Raman Adiabatic Passage (STIRAP)

STImulated Raman Adiabatic Passage (STIRAP) is an adiabatic transfer protocol based on coherent interactions of laser pulses with discrete atomic or molecular systems [39]. STIRAP enables the complete transfer of population between two long lived states, that are optically accessible via a third state. A suitable three level lambda energy level diagram is shown in Fig. 1.3(a). In this system, it is desired to excite molecules from their resting ground state $|1\rangle$ to another state $|3\rangle$. The two states cannot access each other independently. Instead they must rely on an intermediate state $|2\rangle$ to provide this transfer path. The coupling between states $|1\rangle$ and $|2\rangle$ in this system appears as the Rabi frequency $\Omega_1$ and is controlled using a pump laser, whereas the coupling between $|2\rangle$ and the target state $|3\rangle$ is $\Omega_3$ and is controlled by a different laser known as the Stokes pulse. Here, the system can be excited from the initial state $|1\rangle$, to the target state $|3\rangle$, using two delayed Gaussian shaped laser pulses aligned to the so called ‘counter-intuitive’ sequence (Fig. 1.3(b)). During the ‘counter-intuitive’ pulse sequence, the Stokes pulse precedes and overlaps with the pump pulse (Fig. 1.3(b)). Provided that these pulses are varied slowly over time, this results in population transfer that is robust against variations in the Hamiltonian (i.e. changes in the shape of the pulses)[38, 39, 40]. The ordering of the pulses is deemed counter-intuitive as traditionally the pump pulse would precede the Stokes pulse to transfer population from $|1\rangle$ to $|3\rangle$, via $|2\rangle$.

For simplicity, consider all states in the ground state potential, where the Rabi

![Energy Level Diagram](image)

Figure 1.3: (a) A three level energy diagram showing three stable states, (b) the counter-intuitive pulse sequence where $\Omega_3$ precedes the $\Omega_1$ pulse in time, (c) population transfer from $|1\rangle$ to $|3\rangle$ described by the population amplitudes $c_i$ for state $|i\rangle$ when applying the counter-intuitive pulse sequence [39].
frequencies are assumed, $\Omega_i \in \mathbb{R}$. The time varying system Hamiltonian is:

$$H(t) = \hbar \Omega_1(t) (|1\rangle\langle 2| + |2\rangle\langle 1|) + \hbar \Omega_3(t) (|2\rangle\langle 3| + |3\rangle\langle 2|). \quad (1.1)$$

A three state wavefunction can then be defined at any point in time as:

$$|\Psi\rangle = c_1|1\rangle + c_2|2\rangle + c_3|3\rangle, \quad (1.2)$$

where $c_i$ is the probability amplitude of state $|i\rangle$.

Application of the time-dependent Schrödinger equation and using the rotating wave approximation [41] describes the transfer of population between the states with the following coupled differential equations [40]:

$$i \frac{dc_1}{dt} = -\frac{\Omega_1(t)}{2} c_2, \quad (1.3)$$

$$i \frac{dc_2}{dt} = -\frac{\Omega_1(t)}{2} c_1 - \frac{\Omega_3(t)}{2} c_3, \quad (1.4)$$

$$i \frac{dc_3}{dt} = -\frac{\Omega_3(t)}{2} c_2. \quad (1.5)$$

STIRAP is based on the adiabatic ‘dark state’ $|D_0\rangle$, a term that arose due to there being no fluorescence of the intermediate state $|2\rangle$. In the 3-state coupled superposition there also exists odd and even spatially symmetric ‘bright states’ $|D_{\pm}\rangle$. These superpositions, or dressed eigenstates can be written as:

$$|D_{\pm}\rangle = \frac{\Omega_1|1\rangle \pm \sqrt{\Omega_1^2 + \Omega_3^2} |2\rangle + \Omega_3|3\rangle}{\sqrt{2(\Omega_1^2 + \Omega_3^2)}}, \quad (1.6)$$

$$|D_0\rangle = \frac{\Omega_3|1\rangle - \Omega_1|3\rangle}{\sqrt{\Omega_1^2 + \Omega_3^2}}. \quad (1.7)$$

It is important to note that the dark state $|D_0\rangle$ has zero contribution from $|2\rangle$. It is possible to remain in the dark state provided that the coupling terms are varied slowly in time.

Calculating the population transfer through the slowly evolving counter-intuitive pulse sequence is shown in Fig. 1.3(c)[39]. Starting in the ground state $|1\rangle$, small amounts of population are transferred from $|1\rangle$ to $|3\rangle$ until halfway, where there is an even superposition between these two states. Continuing the process shows complete population transfer to the target state $|3\rangle$.

The interesting outcome of STIRAP and the slowly varying counter-intuitive pulse sequence is that it is possible to perform the transfer between states $|1\rangle$ and $|3\rangle$ without the energy ever existing in the intermediate state $|2\rangle$, even though this is the only transfer path possible.
The adiabatic protocol provides a benefit in relative intolerance to timing errors in the pulse sequencing. Population of the intermediate state $|2\rangle$ is ideally suppressed at infinitely long time scales.

The adiabaticity criterion for STIRAP is influenced by how quickly the system Hamiltonian is evolved, and how well maintained the states are throughout population transfer, comparing two states [42];

$$A = \frac{\langle D_0 | \frac{dH}{dt} | D_\pm \rangle}{|\langle D_0 | H | D_0 \rangle - \langle D_\pm | H | D_\pm \rangle|^2},$$

so to satisfy adiabaticity to ensure smooth population transport, $A \ll 1$.

STIRAP has proven useful in molecular dynamics since its inception 25 years ago [38] and has been replicated in many varied fields [39]. Goldner et al. provided the first extension of STIRAP to matter-wave optics [43], with an interest in creating matter-wave beamsplitters. This work provided the first experimental observations of STIRAP-enabled beamsplitters for He-atoms [44] and a matter-wave interferometer for caesium atoms [43].

The first extension of STIRAP outside dynamic laser driven atomic and molecular systems was done by Parkins [45]. Parkins was interested in studying cavity quantum electrodynamics (QED) to manipulate the superposition of a qubit. This cavity-based STIRAP has been demonstrated to generate single photons by passing single cold rubidium atoms through a cavity [46].

The lossless transfer through the intermediate state of STIRAP has its advantages in sensitive quantum information systems. STIRAP has been demonstrated to be useful for efficient qubit manipulation [47] and detection [48] in trapped calcium ions. This robust creation and manipulation of qubit states is paramount in quantum information processing.

A multiple state linkage scheme is important for scalable transport of quantum information as highlighted by Morris and Shore [49]. Adiabatic passage techniques can provide access to multiple states via a chain of states [50, 51]. Even more recently, STIRAP has been proposed to act on an atomic system with N degenerate ground states all connected via a common excited state as described by Fig. 1.4 [52]. This highlights a very powerful feature of this protocol which can theoretically provide population transfer over N states.

The STIRAP protocol serves well at performing robust population transfer between available states and has the potential to couple many states together. This potentially scalable coupling between multiple states is important for exploring higher order quantum interactions. These are two of the properties required to achieve the vision of Section 1.1. However, as these systems provide the transport of energy using resonant Rabi frequencies (i.e. a transfer in the time domain), this is a departure from the spatial transfer required for an integrated photonic
waveguide platform. This particular technique is also not applicable for mass-fabrication, so the search needs to continue to determine if a protocol such as STIRAP is available in a waveguide platform. This would then lead to the development of mass-manufactured waveguide based devices, with the robust transportation capabilities of STIRAP.

1.2.2 Coherent Tunnelling Adiabatic Passage (CTAP)

STIRAP has shown promise in providing some of the criteria required for a quantum information platform, such as the robust transfer of energy and the ability to access many interconnected states. However STIRAP is suited to atomic and molecular systems. An advancement of STIRAP finds a spatial transfer variant of STIRAP, called Coherent Tunnelling Adiabatic Passage (CTAP)[53], enabling population transfer between spatially separated sites.

CTAP was originally envisioned as a method for transporting an electron be-
tween spatially separated quantum wells. A triple donor system is described in Fig. 1.5. Above each donor site sits an ‘S’ (shift) electrode which controls the potential energy of an electron at that site. The overlap between adjacent sites is controlled with the ‘B’ (barrier) electrodes. This wavefunction overlap plays the same role as the laser driven fields of STIRAP. The available quantum dot locations provide the stable storage states and the B electrodes alter the overlap between sites by repelling or attracting the electrons, to provide a counter-intuitive pulse sequence.

The tight binding method was used to approximate the population transfer of the electron between the sites. The Hamiltonian of this 3-state system has previously been described by the coupled differential equations. Equivalently, the Hamiltonian can also be represented in matrix form [53];

$$H(t) = \hbar \begin{bmatrix} 0 & -\Omega_1 & 0 \\ -\Omega_1 & \Delta/\hbar & -\Omega_3 \\ 0 & -\Omega_3 & 0 \end{bmatrix}$$

(1.9)

where $\Omega_1$ and $\Omega_3$ are the tunnelling rates between the central $|2\rangle$, and neighbouring eigenstates $|1\rangle$ and $|3\rangle$ respectively. These coupling terms appeared previously in STIRAP as Rabi frequencies. $\Delta$ is the absolute energy level difference between the central and outer states. The energy of the states can be controlled using the S gates for active compensation. The dressed eigenstates of the Hamiltonian are [53];

$$|D_+\rangle = \sin \Theta_1 \sin \Theta_2 |1\rangle + \cos \Theta_2 |2\rangle + \cos \Theta_1 \sin \Theta_2 |3\rangle$$

(1.10)

$$|D_-\rangle = \sin \Theta_1 \cos \Theta_2 |1\rangle - \sin \Theta_2 |2\rangle + \cos \Theta_1 \cos \Theta_2 |3\rangle$$

(1.11)

$$|D_0\rangle = \cos \Theta_1 |1\rangle + 0|2\rangle - \sin \Theta_1 |3\rangle,$$

(1.12)

with the introduction of mixing terms;

$$\Theta_1 = \tan^{-1}(\Omega_1/\Omega_3),$$

(1.13)

$$\Theta_2 = \frac{1}{2} \tan^{-1}(\sqrt{(2\hbar\Omega_2)^2 + (2\hbar\Omega_3)^2}/\Delta).$$

(1.14)

As mentioned in Section 1.2.1, this adiabatic technique works by maintaining the system in the dark state $|D_0\rangle$ and slowly changing the parameters of $|D_0\rangle$. The dark state is traditionally used to describe an optical phenomenon, however this system can also access an analogous dark state. The coupling terms $\Omega_1$ and $\Omega_3$ follow a similar counter-intuitive pulse sequence to STIRAP, except now the pulse will be electrical, rather than laser driven.
The energies of the dressed eigenstates are [53]:

\[ \epsilon_{\pm} = \frac{\Delta}{2} \pm \frac{1}{2} \sqrt{(2\hbar \Omega_1)^2/(2\hbar \Omega_3)^2 + \Delta^2}, \]

(1.15)

\[ \epsilon_0 = 0. \]

(1.16)

Fig. 1.6 shows the eigenenergies of the three dressed states \( |D_+ \rangle \), \( |D_- \rangle \) and the previously introduced ‘dark state’ \( |D_0 \rangle \), which has no contribution from site \( |2 \rangle \) as described in (1.12). The maximum split in energy levels observed at the middle point of the protocol is consistent with the electrodes providing maximum coupling. The middle position also marks where the adiabaticity parameter is the largest and is therefore the least adiabatic. With an adequate time scale the dark state \( |D_0 \rangle \) is isolated from the bright states and will provide a robust transport path for the electron.

This particular study considers ideal Gaussians to alter the coupling strength for an appropriate counter-intuitive, or intuitive sequence as shown in Fig. 1.7(a) [53]. As this is an adiabatic technique the system must be evolved slowly in order to achieve robust transportation. This is verified by changing the operational time scale of the gate and monitoring the final state population of the target state \( |3 \rangle \) (Fig. 1.7(b)). As the time scale increases so does the fidelity. The adiabaticity criteria can again be determined using equation 1.8, by ensuring \( A \ll 1 \).

Fig. 1.7(c) shows successful spatial adiabatic transport between states \( |1 \rangle \) and \( |3 \rangle \) as the system is slowly evolved through the counter-intuitive sequence. There is relatively no interaction at the start, when both couplings are small. At half way through the system evolution there is now an even amount of electron distribution between \( |1 \rangle \) and \( |3 \rangle \). Continuing through the protocol shows the complete transfer to \( |3 \rangle \). As predicted by STIRAP, provided adiabaticity is satisfied, the mediating state \( |2 \rangle \) is not populated and the transfer is lossless. The intuitive sequence provides a reversed sequence of operation (Fig. 1.7(c)). Evolving the system with this sequence results in a sinusoidal output dependence with increased total time (Fig. 1.7(e)), and strong population and interaction with the intermediate state.
Figure 1.7: (a) Counter-intuitive coupling scheme turns on gate $B_{23}$ before $B_{12}$, (b) population in site $|3\rangle$ at the end of the transport for varying time scales, (c) population transfer between all three states for the counter-intuitive sequence, (d) swapping the sequence of the gates provides the intuitive scheme, (e) final population of site $|3\rangle$ for the intuitive sequence at increasing time scales, (f) population levels of all three states for a particular time scale that provides transfer to $|3\rangle$ using the intuitive sequence. Adapted from [54].

This is an important result, as it highlights that the adiabatic protocol introduced by STIRAP can be applied to conserving and transferring quantum states over some distance. The discussed particle based CTAP devices extends STIRAP to spatially dislocated states and like STIRAP can provide robust information transfer between these well defined states. However, this research is interested in performing this energy transfer within a waveguide platform.

1.2.3 Waveguide CTAP

CTAP has recently been applied to an optical platform, demonstrating robust light transfer among evanescently coupled optical waveguides [40, 55, 56]. These designs consisted of a geometrically curved three waveguide directional coupler (Fig. 1.8(a)) to provide the conditions for adiabatic passage [56, 55]. Longhi later fabricated this design to show CTAP in Ag-Na diffused optical waveguides [40, 57].
This system is designed to use three waveguides ($|1\rangle, |2\rangle, |3\rangle$) that can act as spatially separated states (Fig. 1.8(b)), analogous to the energy levels of STIRAP (Fig. 1.3(a)). The tunnel matrix elements for the waveguides are the proximal evanescent coupling coefficients between the centre and either neighbouring waveguide ($\Omega_1$ and $\Omega_3$). The curved trajectory provides a suitable counter-intuitive coupling sequence as a function of the forward propagation direction ($z$) (Fig. 1.8(c)). Both curves are offset a distance ($\delta$) in the $z$ direction. As the coupling strength is an exponential function, this geometry will provide two overlapping Gaussians, shifted from each other by a distance $\delta$, playing the role of the delay between the Stokes and pump laser pulses in STIRAP.

The goal here is again to counter-intuitively transfer light completely from $|1\rangle$ to $|3\rangle$. It is important to note that $|1\rangle$ and $|3\rangle$ cannot communicate independently and can only be connected by using $|2\rangle$ as a mediating nearest neighbour. Simulations have been provided by solving the coupled mode equations of the Schrödinger like paraxial wave equation [40] and is again found to be analogous to the STIRAP equations.

As light is launched into $|1\rangle$, it will stay guided as the two waveguides approach $|2\rangle$ (Fig. 1.9(a-b)). At the beginning of the protocol there is little obvious light interaction with $|2\rangle$ and $|3\rangle$ even as $|3\rangle$ continues to approach $|2\rangle$. At the end of the protocol all the light exits waveguide $|3\rangle$.

Exciting waveguide $|3\rangle$ shows a very different result (Fig. 1.9(c-d)), light will initially be completely guided by $|3\rangle$. As it approaches $|2\rangle$, light directionally couples between $|3\rangle$ and $|2\rangle$ with a frequency that is dependent on how close the waveguides are to each other. As the light continues to propagate in the $z$ direction waveguide $|1\rangle$ is now close enough to interact with $|2\rangle$ and at some point light will hop between all three waveguides. Then, as the coupling to $|3\rangle$ is reduced, light no longer interacts with waveguide $|3\rangle$, and $|1\rangle$ and $|2\rangle$ continue to beat.
nally, \( |1\rangle \) deviates far enough away to no longer couple with \( |2\rangle \) and there will be a signal output from all three waveguides that depends on all of those interactions. The use of directional coupler beating makes this device very sensitive to fabrication tolerances in every dimension and highlights the robustness offered with the counter-intuitive sequence.

The transfer of light between waveguides using coupled mode equations has been proven to be consistent with the STIRAP equations (1.3-1.4) and the previously discussed quantum dot CTAP evolution of the Hamiltonian [40].

The small field present in \( |2\rangle \) during the counter-intuitive sequence is due to a number of reasons; such as a non-optimised coupling scheme, and these have been explored at length in various studies [58, 59]. Also, as this is an adiabatic protocol it can only promise ‘perfect’ transport to the target waveguide as the Hamiltonian is evolved infinitely slowly, however the dramatically different population transfer behaviour observed for the counter-intuitive case is proof that adiabatic passage is possible within a finite lengthed device.

This new branch of adiabatic waveguide devices seem promising in performing robust transfer of photons between waveguides in an integrated circuit. The current devices shown, however are limited to nearest neighbour couplers. A third integrated waveguide has been shown to provide another spatial dimension [60, 61] and other groups have attempted to use femto-second laser direct-write techniques to provide this third dimension of spatial control in projects unrelated to CTAP, but related to quantum information systems, allowing three waveguides to communicate with the same magnitude [30, 31]. As described earlier the more

Figure 1.9: Experimental and simulated light population transfer within a three waveguide CTAP directional coupler: (a) Light input into the counter-intuitive input, waveguide \( |1\rangle \), (b) counter-intuitive simulation, (c) exciting waveguide \( |3\rangle \) provides the intuitive response, (d) intuitive simulations. From [40].
states that can be linked together simultaneously, the higher dimensional Hilbert space that can be operated on. However, this still only allows a $3 \times 3$ device, not an arbitrary $N \times N$. The 3D fabrication technique is also not suitable for mass-manufacture in a silicon foundry setting which is confined to planar topologies.

1.2.4 Summary

A thorough analysis of an adiabatic transfer protocol shows that it has advantages in providing well controlled robust signal transfer. Such a feature is important for highly sensitive photonic experiments. The described adiabatic three waveguide device can transport light across an intermediate waveguide. However, this intermediate waveguide is needed to facilitate the transfer. As the coupling mechanism between each waveguide is based on nearest neighbour interactions, SWAP gates are required to transport light across large distances. Adiabatic devices are larger than their non-adiabatic counterparts so device footprint becomes an issue with each additional operation. To solve the problem of nearest neighbour interactions, the answer must lay outside of STIRAP and CTAP.

1.3 Lateral leakage of shallow ridge silicon waveguides

The discussion of adiabatic passage techniques in Section 1.2.1 provided a robust and well controlled transfer path for photons which is of interest in sensitive quantum interferometry and entanglement operations. The developed waveguide structures require nearest neighbour interactions to operate, which severely limits the complexity of allowable devices.

In 2007, Webster et al. fabricated a series of silicon-on-insulator (SOI) waveguides of varying widths and lengths [62]. The waveguides were designed with a shallow etched ridge to reduce scattering losses due to unideal sidewall roughness. These waveguides performed very well with extremely low loss for the guided transverse-electric (TE) mode. However, surprisingly the transverse-magnetic (TM) mode experienced severe losses, for all but two of the waveguide widths.

This effect was first predicted by Peng and Oliner in 1978, for strip and rib waveguides, but conceived at RF frequencies [63]. They also predicted that the leakage effect could be cancelled with suitable waveguide engineering.

Due to the high index contrast between the air cladding ($n=1$) and silicon core ($n=3.5$) the TE and TM guided modes are not ‘pure’. This results in the primarily vertically oriented TM-like mode having a non-zero longitudinal field component. There also exists vertically confined, but laterally unguided TE slab modes. These
TE slab modes can have a longitudinal field component if they propagate at an angle to the propagation axis. This mode hybridisation provides a non-zero overlap between TM and TE slab modes, which can allow them to couple.

The effective indices ($n_{\text{eff}}$) or propagation velocity ($\beta$) of the guided TM and TE modes are far too different to allow phase matching and are mutually isolated. However, the guided TM mode can be phase matched to a higher effective index TE ‘slab’ mode, provided that it propagates at an angle ($\theta_{\text{clad}}$) to the waveguide (Fig 1.10(a)). Where, $\theta_{\text{clad}}$ can be calculated as [64]:

$$\theta_{\text{clad}} = \cos^{-1}(N_{\text{eff,TM}}/n_{\text{TE clad}}).$$  (1.17)

where $N_{\text{eff,TM}}$ is the effective refractive index of the TM guided mode and $n_{\text{TE clad}}$ is the effective refractive index of the TE slab mode outside of the core region.

The waveguide sidewall interface provides the perturbation necessary for these two modes to couple (Fig. 1.10(b))[65, 66]. The consequence of this TM-TE mode coupling is highlighted in Fig 1.10(c) which shows a cross-sectional view of a standard SOI shallow ridge waveguide with representations of the field presence of the guided but hybridised TM mode (blue) and the vertically bound, but

Figure 1.10: The fundamental TM mode is altered and hybridised in shallow ridge waveguides: (a) At certain in-plane angles, the radiative TE modes of the silicon slabs ($\beta_{\text{TE core}}$ and $\beta_{\text{TE clad}}$) can be phase matched to the fundamental guided TM mode ($\beta_{\text{TM}}$), (b) interference is responsible for the appearance of ‘magic’ values where the radiation produced by either sidewall is interfering destructively, and this effect is related to the width ($w$) of the waveguide and the operational wavelength ($\lambda$), (c) a representation of the mode profile of the hybridised TM-like mode containing both vertical and horizontal components.
laterally free TE slab mode (green). If the coupling between these two modes is strong, there will be inherently large losses experienced for the guided TM mode. Fortunately, this effect can be mitigated by choosing the correct boundary conditions. As the TM ray is incident on a waveguide boundary it generates a reflected \((T_{E_{core}})\) and transmitted \((T_{E_{clad}})\) TE rays due to the TM-TE mode coupling (Fig. 1.10(b)). This occurs at both waveguide sidewalls, producing four TE rays that can interfere with one another. If the rays traversing the waveguide core arrive out of phase when entering the cladding region, the TE radiation will be coherently cancelled and is referred to as ‘magic width’. If instead the rays arrive in phase with the \(T_{E_{clad}}\), this will lead to extremely high TM/TE coupling as the guided light ‘laterally leaks’ into the cladding.

The resonant ‘magic’ width \((W)\) can be calculated using \([62, 65]\):

\[
W = \frac{(m + \Delta \phi/(2\pi))\lambda}{\sqrt{(n_{\text{eff,TE}}^{\text{core}})^2 - \lambda_{\text{eff,TM}}^2}}
\]

(1.18)

where \(m\) is an integer, \(\Delta \phi\) is the phase difference between the reflected and transmitted TE wave at the ridge interface, \(\lambda\) is the operating wavelength and \(n_{\text{eff,TE}}^{\text{core}}\) is the effective index of the TE slab mode within the core region.

Simulations of the lateral leakage mechanism require a fully vectorial mode solver with appropriate boundary conditions \([65]\). The ridge width dependence on the lateral leakage loss for a standard shallow ridge SOI foundry ready waveguide (Fig. 1.11(a) inset) is simulated using fully open lateral boundaries. With open

Figure 1.11: (a) The width dependence of the propagation loss of the fundamental TM mode, (b) at ‘magic width’ \((w = 722 \text{ nm})\) the lateral radiation cancels out and the waveguide has a low propagation loss. (c) at ‘anti-magic width’ \((w = 1.1 \mu\text{m})\) the radiation is reinforced and results in a strong field presence in the cladding.
boundaries, the imaginary component of the TM guided mode effective index represents the magnitude of the lateral leakage loss. The width of the waveguide is adjusted from 400 nm to 1500 nm and the imaginary effective index recorded and converted to dB/mm (Fig. 1.11(a)). At the smallest waveguide width, there is a very strong leakage interaction, this slowly decreases until a sharp resonance is seen for a width of 722 nm. This is the value of the first ‘magic width’ for this particular geometry. Increasing the waveguide width further starts radiating more power, until it reaches a maximum (‘anti-magic’ width). The waveguide loss again decreases until the next resonant lossless width is reached. The lateral field plots of the fundamental TM guided mode for ‘magic’ and ‘anti-magic’ widths in a single shallow ridge are shown in Fig. 1.11(b-c). The annotations serve to highlight the fact that, if the radiation produced by each sidewall accrues a phase change of $\pi$ as it arrives at the second waveguide boundary, then the radiation will be coherently cancelled and the TM guided mode will provide extremely low losses. Conversely for ‘anti-magic’ width, both sidewall produce radiation that sums and produces large losses.

1.3.1 Lateral leakage in coupled structures

As additional shallow ridge waveguides are introduced to a common silicon slab, they are highly likely to interact with each other. Consider two, non-‘magic’ width shallow ridge waveguides on a common slab as per the inset of Fig. 1.12(a) with increasing waveguide separation [67].
At 2 \( \mu \text{m} \) separation the odd and even supermodes have very different refractive indices as expected by evanescent, nearest neighbour coupling (Fig. 1.12(a)) [67]. As the waveguide separation increases to 5 \( \mu \text{m} \) there is no longer any evidence of evanescent coupling. Instead of decreasing to a minimum, the supermode effective indices oscillate. The waveguides continue to act at a distance that is now sinusoidal rather than exponential. As the real part of the effective indices of the two supermodes oscillates, the imaginary component oscillates out of phase with the real components (Fig. 1.12(b)). Observation of the laterally polarised mode fields of the two guided TM supermodes (Fig. 1.12(c-d)) shows that these two waveguides are inherently linked with each other and the TE slab mode. This study was extended further, highlighting how it is possible to launch a coherent TE beam with a single lateral leakage waveguide [64]. If this light can be collected by another lateral leakage waveguide, these two will act similarly to an antenna transmitter, receiver pair. This type of behaviour suggests that this platform may be suitable for using this TE slab radiation as an optical bus.

1.3.2 Summary

This section highlighted a mechanism that can provide long-range coupling between shallow ridge silicon photonic waveguides. This could provide the basis of an optical bus, where the participation of each waveguide can be selectively controlled with waveguide engineering. The adiabatic techniques highlighted in Section 1.2.1 could be coupled with this long-range optical bus, to provide robust transfer between waveguides, without exciting the bus mode itself.

1.4 Thesis outline

This thesis presents a new technique for transporting light between integrated photonic waveguides via a bus, that no longer relies on nearest neighbour interactions. The bus operation is suitable for use within mass-manufacturable industry standard shallow etched silicon waveguides. This platform has been conceived as a means to achieve highly robust, multi-dimensional quantum gates.

Chapter 2 investigates the feasibility of reliably fabricating lateral leakage waveguides in a standard silicon photonics platform. The devices are simulated using eigenmode expansion techniques [65] and the waveguides are designed using IPKISS [68]. Fabrication took place through IMEC as part of a multi project wafer (MPW) run, providing substantial cost reductions to the project. A variety of different waveguides are designed and fabricated on a standard 220 nm thick, 70 nm etched, air-clad silicon-on-insulator platform [69]. A full experimental analysis is provided, with the results agreeing very well with predictions.
This study shows that waveguides exhibiting lateral leakage can be developed in a mass-fabricable silicon platform. A method for electronically tuning the leakage resonances using an active liquid crystal overlayer is also reported. This improves the confidence in developing more complicated devices that make use of the TM guided mode in shallow etched waveguides.

Chapter 3 proposes a technique suitable for transferring light between vastly separated waveguides using the lateral leakage mechanism introduced in Chapter 2. Control over the lateral leakage of multiple waveguides on a common slab has lead to the development of an all photonic bus. Full vector mode matching techniques are employed to simulate the transfer of light between waveguides and this new optical bus. Combining the lateral leakage coupling mechanism with a robust adiabatic technique (CTAP) [53] results in a large reduction in the amount of power lost due to exciting the optical bus. In fact for ideal operation, there is absolutely no excitation of the optical bus whatsoever. This technique exhibits low loss and robust light transfer between waveguides separated well above the limit for evanescent coupling. The adiabatic technique also provides added protection against variations in absolute device length and precise coupling conditions when compared to traditional directional couplers, which require tunable heater elements. This development breaks the topological restraints of waveguide positioning currently imposed by traditional methods of nearest neighbour interactions and can lead to more exotic device design.

In Chapter 4, the focus of the newly developed adiabatic optical bus mode shifts from photonic circuit engineering to an application in quantum photonics. This study suggests that careful control over the lateral leakage coupling to each waveguide can incur a phase rotation on entangled photon states. A variety of devices are proposed and simulated using tight-binding conditions. This study lays the ground work for the development of fundamental building blocks of a linear optics quantum computer, namely controllable adiabatic multi-input/output beam splitters and demonstrates functions required for an arbitrary linear optics quantum computer such as $\frac{1}{3}:\frac{2}{3}$ splitting, Hadamard operation ($\frac{1}{2}:\frac{1}{2}$) and controlled NOT gating.

The final chapter concludes on the major findings of this thesis and provides an outlook towards future research endeavours.

The primary objectives of this research is in the development of a new waveguide platform. The important feature of this new platform is the long-range, selective communication between waveguides that is no longer solely restricted to nearest neighbour interactions. The waveguides can communicate instead via an optical bus. Breaking this restriction, allows waveguides to communicate in a manner that is simply not possible in current planar photonic platforms. The optical bus introduced in this work is a missing link in photonic circuits, helping to realise highly complex, yet scalable integrated photonic systems such as,
high-order quantum interferometry, or a complex photonic routing network. The application of an adiabatic passage protocol increases the fidelity of this transport, and reduces cross-talk between neighbouring waveguides making it suited towards sensitive photonic systems. The described silicon photonic long-range coupling platform is unprecedented and will allow a new suite of scalable classical and quantum photonic devices.
1.5 Publications and conferences originating from this doctorate

International Conferences


National Conferences


Peer Reviewed Journal Articles


Peer Reviewed Journal Article - Second-Author

Chapter 2

Standard silicon foundry fabricated shallow ridge waveguides

2.1 Quantitative analysis of TM lateral leakage in foundry fabricated silicon rib waveguides

Quantitative analysis of TM lateral leakage in foundry fabricated silicon rib waveguides

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Abstract—We show that thin, shallow ridge, silicon-on-insulator waveguides exhibiting lateral leakage behaviour can be designed and fabricated using a standard silicon photonic foundry platform. We analyse the propagation loss through observation of the transmitted TM polarized mode and TE polarized radiation and experimentally demonstrate that propagation losses as low as 0.087 dB/mm can be achieved. This demonstration will open a new frontier for practical devices exploiting lateral leakage behaviour with potential applications in the fields of biosensing and quantum optics among others.

I. INTRODUCTION

It has been shown that silicon waveguides with exceptionally low loss can be achieved using a shallow ridge configuration [1]. It was predicted that these waveguides should exhibit strong vertical evanescence when operated in the TM polarization, with potential applications for hybrid integration [2] and biosensing [3]. However, it was found experimentally that the TM mode only achieved low loss transmission for a discrete set of waveguide widths, an effect not predicted by traditional simulators used for initial predictions [4].

This phenomenon has been theoretically explained due to phase matching (Fig. 1(a)) of the fundamental guided TM mode of the ridge and the TE radiation (in the cladding) [4], [5]. The magnitude of this coupling is controlled by interference effects (Figs. 1(b,c)). Certain waveguide widths either coherently cancel the radiation (Figs. 1(d,e)), or completely reinforce it resulting in the strongest losses (Fig. 1(f)).

It follows naturally that this strong polarization dependent radiation can be used to realize compact integrated polarizers [6] or variable attenuators [7], but rather than viewing this effect as a simple loss mechanism, one can consider it as a means of generating controlled coherent beams of TE light [8], with one possible application being long-range transfer of an optical carrier between apparently uncoupled waveguides [9]. The analysis of most of the proposed structures harnessing lateral leakage has been limited to simulation. Implementing the experimental approach of [1], [4] requires non-standard processing which may be difficult to access.

In this paper, we show that lateral leakage can be reliably achieved using standard silicon foundry fabrication. We observe the width dependent leakage loss through the transmission of the TM mode and also observe the leaked TE radiation directly. Our observations match well with predictions. This standard fabrication platform offers benefits such as high precision, repeatability, throughput and the possibility of inclusion in multiple project wafer initiatives (MPW), lowering fabrication costs. This demonstration will improve the understanding and confidence in the operation of lateral leakage so that new and interesting devices can be developed and adopted into industrial integrated silicon photonic systems. Preliminary material from this work was published in [10]; however, here we present deeper analysis and substantial additional results.

![Fig. 1: The fundamental TM mode is altered and hybridized in shallow etched waveguides; (a) At certain in-plane angles, the radiative TE modes of the silicon slab (β_{TE,1,0} and β_{TE,2,0}) can be phase matched to the fundamental guided TM mode (β_{TM}). (b) Interference is responsible for the appearance of ‘magic’ values where the radiation produced by either sidewall is interfering destructively, and this effect is related to the width (w) of the waveguide and the operational wavelength (λ), (c) a representation of the mode profile of the hybridized TM-like mode containing both vertical and horizontal components, (d) the width dependence of the propagation loss of the fundamental TM mode, (e) at ‘magic width’ (w = 722 nm) the lateral radiation cancels out and the waveguide has a low propagation loss, (f) at ‘anti-magic width’ (w = 1.1 μm) the radiation is reinforced and results in a strong field presence in the cladding.](image-url)
II. DESIGN OF THIN SHALLOW RIDGES FOR CMOS FABRICATION AND GRATING COUPLER CHARACTERISATION

This section aims to show that thin shallow ridge waveguides exhibiting lateral leakage can be realized using standard silicon photonic fabrication. We chose the ePIXfab IMEC standard platform using 220 nm thick silicon on 2 μm silicon dioxide, patterned with either 70 nm or full 220 nm etching with minimum features sizes of 150 and 130 nm, respectively [11]. The upper cladding was air. The devices were designed and simulated using the ‘IPKISS’ framework [12] and a custom eigenmode expansion model [13] to enable simulation of the lateral leakage behaviour. Our simulations predicted that ridges with these parameters, should support a single TM mode at a wavelength of 1550 nm up to a width of about 850 nm. Fig. 1(d) presents the predicted lateral leakage loss as a function of ridge width, clearly showing low-loss ‘magic’ width behaviour at a width of 722 nm.

We designed a series of waveguides with a range of widths from 400–800 nm and lengths 10 μm–1 mm. Each waveguide was interfaced to single mode fibers using tapered focussed TM grating couplers [14], [15] designed for optimum efficiency in the λ=1500–1570 nm range and phase matched only to the TM mode. To observe the TE lateral leakage radiation, a series of grating couplers, phase matched and oriented to align with the expected TE radiation, were included on either side of the waveguide along its length. The very different phase velocities of the TE and TM slab modes ensures only the TE mode is out-coupled by this grating. Fig. 2(a) presents a schematic of the designed structure and Figs. 2(b-c) illustrate the methods for measuring the TM waveguide transmission and the TE lateral leakage radiation, respectively.

The designed devices were fabricated, and to determine the actual waveguide dimensions, a sample of devices were cross-sectioned and imaged using scanning electron microscopy (SEM). An SEM image of one of the realized structures is presented in Fig. 3. The silicon thickness was found to be 215 nm, the ridge was 70 nm and the ridge widths tended to be

10–20 nm wider than the nominal design.

Simulating with the slab thickness at 215 nm, we found that the predicted ‘magic’ width is 717 nm at a wavelength of 1550 nm. From this point, all reported dimensions are as-fabricated.

To qualitatively test the structures, an infrared camera was positioned above the 0.5 mm long waveguides and the input coupler was excited with a wavelength of 1550 nm. Fig. 4(a) shows a visible wavelength microscope image of one of the waveguide structures. Fig. 4(b) shows the infrared image of the 720 nm wide waveguide excited at 1550 nm. On the input side we see scattered light. On the output side we see a very bright spot corresponding to light exiting through the output TM grating coupler. In between we see darkness indicating that light is confined within the waveguide.

Fig. 4(c) shows an infrared camera image of the 675 nm wide waveguide excited at 1550 nm. Again we see scattered light from the input. The light at the output TM grating coupler is dimmer than in Fig. 4(c). In between we see excitation of the TE side gratings, decreasing with distance from the input. Fig. 4(c) presents clear evidence of lateral leakage behaviour in thin-ridge waveguides while Fig. 4(b) presents clear evidence of the suppression of lateral leakage when the waveguide is ‘magic’ width. Together these represent the first evidence of lateral leakage observed in mass fabricated silicon waveguides.
III. WAVEGUIDE CHARACTERIZATION

A. Width dependent propagation loss

We first quantify the propagation loss as a function of waveguide width using the cut-back method, measuring the transmission efficiency for identical waveguides of different lengths and then extract the loss per unit length. The transmission of waveguides of widths 420, 720, 760 and 820 nm and lengths between 10 and 1000 µm were each measured at a wavelength of 1550 nm and the transmission efficiency as a function of length is plotted in Fig. 5. To eliminate the response of the grating couplers, a fully etched wire waveguide of width 420 nm with identical grating couplers was used as a normalization reference. The propagation loss (α) is the slope of the line in dB/mm. These results clearly show that there is a strong dependence of observed loss on waveguide width with α=0.357 dB/mm when w=720 nm rising to α= 585.95 dB/mm, similar to the behaviour reported in [4].

B. Wavelength dependent propagation loss

Lateral leakage is a resonant effect and thus we expect the loss to not only depend on ridge width, but also on wavelength [4]. For 1550 nm, we expect ‘magic’ width behaviour, where the lateral leakage loss is completely cancelled, at a width of 717 nm. To characterize the wavelength dependent propagation loss, we measured the power transmitted through waveguides with widths of 675, 700, 720 and 740 nm over the entire fabricated length range, while scanning the wavelength from 1500–1570 nm as illustrated in Fig. 2(b). The loss of each waveguide width was again determined using a linear fit.

Fig. 6 presents the transmitted power per unit length for each waveguide as a function of wavelength. Resonant loss cancellation is clearly evident with waveguide widths of 700 nm and 720 nm exhibiting minimal loss of 0.238 dB/mm and 0.087 dB/mm at wavelengths of 1500 nm and 1563 nm, respectively.

The expected wavelength dependent lateral leakage behaviour was simulated using eigenmode expansion [13] and these predictions are also presented in Fig. 6. An excellent match to the measured values is evident, particularly the 720 nm width waveguide with transmission loss approaching zero at the resonant wavelength.

C. Conversion of TM guided propagation to TE radiation

To observe the TE radiation we use side gratings, with period and orientation for efficient diffraction of the TE mode propagating at the expected angle of 45 degrees, as illustrated in Fig. 2(c). Waveguides of fabricated widths w = 675, 720 and 810 nm and length 500 µm were excited with λ=1550 nm and the intensity measured at the TE side gratings as a function of length is presented in Fig. 7. The rapid oscillations are due to the numerous, discrete side gratings. A linear fit to each trace indicates exponential decay with the slope yielding the attenuation coefficient. These are similar to those of Fig. 6 indicating that the observed width dependent loss is due to conversion of the TM guided light into lateral TE radiation.

Fig. 5: TM transmission at λ=1550 nm vs waveguide length for different waveguide widths. The gradient yields the waveguide propagation loss (α).

Fig. 6: The wavelength dependant TM transmission loss of four fabricated waveguide widths. Simulations have been performed to compare the spectral response of each waveguide width.

Fig. 7: Scanning an output fiber across the center of these side gratings provides quantitative information regarding the TE conversion and the loss can be estimated with a linear fit.
Fig. 8: (a) Experimental propagation loss obtained from the transmission measurements and side grating measurements, compared to the simulated loss at \( \lambda = 1550 \text{ nm} \). (b) Simulation of the propagation loss (\( \alpha \)) expected for 215 nm high, 70 nm etched air clad SOI waveguides as a function of waveguide width and wavelength, the ideal magic width is dashed white. (c) Propagation loss (\( \alpha \)) calculated from the interpolated experimental transmission data, which is in good agreement with the simulation.

IV. ANALYSIS OF LEAKAGE MEASUREMENTS

We have characterized the fabricated thin ridge waveguides over a wide range of wavelengths, widths and lengths, and we have compared these measurements with simulations based on the measured as-fabricated waveguide parameters. Cut-back analysis of the direct TM transmission loss (Fig. 5) and the observation of the TE radiation (Fig. 7) was used to extract the propagation loss. Fig. 8(a) presents a survey of the propagation loss as a function of waveguide width at \( \lambda = 1550 \text{ nm} \) extracted from both the measured TM transmission and TE radiation. The vertical and horizontal error bars represent uncertainty in the measured loss and dimensional accuracy (\( \pm 10 \text{ nm} \)), respectively. The predicted loss of the fundamental waveguide mode due solely to TE radiation, is also presented in Fig. 8(a). Excellent agreement between both measurements and prediction is evident up to widths of 760 nm. At widths above 780 nm, the waveguide is expected to begin to support a strongly radiating higher order resonance, which becomes a guided mode at 850 nm, and this may explain the additional transmission loss observed at these widths and above.

A wavelength sweep was also conducted during the measurements. Fig. 8(b) and (c) present the predicted and interpolated, experimentally measured propagation loss as a function of waveguide width and \( \lambda \), respectively. Again we find a very good match, with a slight skew towards the higher waveguide widths.

V. CONCLUSION

We have shown that thin, shallow-etched SOI waveguides exhibiting lateral leakage behaviour can be realized using standard, silicon photonic foundry fabrication. We have characterized the waveguides via analysis of the transmitted TM mode and also direct observation of the TE radiation. We find very good correspondence between eigenmode expansion predictions and two different measurement procedures. Magic width behaviour with resonant cancellation of the leakage leading to losses as low as 0.087 dB/mm have been observed experimentally. The ability to design, simulate and experimentally realize these passive waveguide structures using a standard silicon photonic fabrication platform creates new opportunities for practical design and accessible fabrication of realizable devices exploiting lateral leakage with potential applications including biophotonic sensing and quantum information processing.

REFERENCES


2.2 Electrically tunable lateral leakage loss in liquid crystal clad shallow-etched silicon waveguides

Electrically tuneable lateral leakage loss in liquid crystal clad shallow-etched silicon waveguides

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Abstract: We demonstrate electrical tuning of the lateral leakage loss of TM-like modes in nematic liquid crystal (LC) clad shallow-etched Silicon-on-Insulator (SOI) waveguides. The refractive index of the LC layer can be modulated by applying a voltage over it. This results in a modulation of the effective index of the SOI waveguide modes. Since the leakage loss is linked to these effective indices, tunable leakage loss of the waveguides is achieved. We switch the wavelength at which the minimum in leakage loss occurs by 39.5nm (from 1564nm to 1524.5nm) in a 785nm wide waveguide. We show that the leakage loss in this waveguide can either be increased or decreased by modulating the refractive index of the LC cladding at a fixed wavelength.

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OCIS codes: (200.4650) Optical interconnects; (230.3720) Liquid-crystal devices; (230.7390) Waveguides, planar; (230.7405) Wavelength conversion devices; (250.5300) Photonic integrated circuits; (250.6715) Switching.

References and links
1. Introduction

In recent years, Silicon-on-Insulator (SOI) photonics has successfully established itself as a viable technology for photonic integrated circuits, especially for optical interconnect applications [1, 2]. This success is in large part due to intrinsic advantages which arise as a result of working within the SOI material platform: the high index contrast makes it feasible to design waveguides with very narrow bends, making high density integration a practical reality. In addition to this, compatibility with existing complementary metal-oxide semiconductor (CMOS) fabrication facilities means SOI photonics isn’t impeded by expensive start up costs. The two main types of waveguides used in SOI technology are deep-etched (strip/wire) and shallow-etched (ridge) waveguides. Propagation losses of 0.27dB/cm and 1 to 2dB/cm have been reported for shallow-etched and deep-etched waveguides respectively [3, 4]. This indicates that shallow-etched waveguides are suitable for applications where an optical signal needs to be efficiently transmitted over a distance of a few tens of centimeters e.g. on-chip optical interconnects. Due to their geometry, shallow-etched waveguides also allow lateral electrical access; this makes them suitable for active devices like waveguide modulators and lasers.

Because of their high index contrast, silicon waveguides are extremely birefringent: the TE-like mode and the TM-like mode have very different propagation constants. Therefore, silicon waveguides are usually used only for the TE polarization. Strip waveguides also guide TM-polarized light, but in shallow-etched waveguides, it has been theoretically and experimentally demonstrated that the TM-like mode suffers from waveguide-width-dependent lateral leakage loss [5, 6]. The effective index of the guided TM-like mode in such waveguides is comparable to that of the radiating (cladding) slab TE-like mode. This radiating slab TE-like mode can propagate in any direction since it is unguided and can be phase matched to the guided TM-like mode at a particular angle. The guided TM mode thus suffers from lateral leakage loss since it can be phase matched to a radiating mode. Accordingly, the TM-like mode in a shallow-etched waveguide is always lossy except for the case when the waveguide satisfies the well known resonance condition [5].

\[
W = \frac{m\lambda}{\sqrt{n_{\text{eff,TE}}^2 - n_{\text{eff,TM}}^2}}
\]  

(1)
Where \( W \) is the waveguide width, \( \lambda \) is the wavelength of the light in vacuum, \( m \) is a positive integer denoting the order of the resonance condition, \( N_{\text{eff, TM}} \) is the effective index of the guided TM-like waveguide mode and \( n_{\text{eff, TE}}^{(\text{core})} \) is the effective index of the (unguided) TE wave which traverses the waveguide core. Conversely, for a waveguide with a fixed width, there exists a wavelength at which the leakage loss is minimal.

In addition to straight waveguides, this lateral leakage behavior has been reported to be present in shallow-etched bent waveguides and ring resonators [7] as well. Several attempts have been made to mitigate this undesirable effect [8–10]. In the current work, we propose a method for actively tuning the position of the leakage loss minima in liquid crystal (LC) clad shallow-etched waveguides by applying a voltage over the LC.

The silicon devices were designed using the IPKISS parametric design framework and fabricated by IMEC through the ePIXfab multi-project-wafer service. They consist of a 220nm silicon layer on a 2\( \mu \)m buried oxide layer. The grating couplers and the shallow-etched waveguides are defined by a 70nm etch into the silicon, while the other waveguide structures (such as the grating coupler tapers) are fully etched. For our experiments we did not have an oxide cladding deposited, so the top silicon surface is exposed to air. Figure 1 shows a sketch of the geometry of the waveguide and the coordinate system we use.

![Fig. 1. Geometry of a shallow-etched SOI waveguide. The z-axis is perpendicular to the plane of the paper.](image)

We design the waveguides to be 1 cm long with varying widths. The widths of the waveguides are chosen so that they are close to the first leakage loss minimum in the desired wavelength range as expected from previous theoretical considerations [5, 7, 11, 12]. The waveguides have grating couplers placed at their extremities. The grating couplers are optimized for TM polarized light. They are curved and have a period of 1050nm with a fill factor of 50%. Light is coupled into and out of the waveguides by placing optical fibers inclined at an angle of 10° with the vertical above the grating couplers.

The remainder of the paper is organized as follows. In Section 2, we present measurements of the leakage loss of air-clad waveguides and compare them to previous results reported in the literature. Section 3 is devoted to the study of LC clad waveguides. Here we give all details concerning the assembly of the LC cladding over the shallow-etched waveguides. Section 4 deals with a discussion of the phenomenology uncovered by the experiments we perform. We conclude the paper in Section 5.

### 2. Air cladding

The purpose of the air-clad measurements is two-fold. First, we need to ensure that the waveguides we work with exhibit the leakage loss behavior as expected. Second, we intend to study the potentially complicated effect a LC cladding has on the leakage properties of a waveguide. Accordingly, a measurement of the air-clad case provides a bench mark for the more complicated case. We focus on the wavelength window from 1510nm to 1590nm which is determined by the experimental setup.

A loss measurement is carried out as follows; we use a fiber-coupled tuneable laser to couple light into the shallow-etched waveguide through one grating coupler. The light is coupled out of
the waveguide at the other grating coupler into a fiber-coupled power meter. Each measurement is corrected for the contribution of the grating couplers by subtracting the transmission of a short deep-etched waveguide with identical grating couplers at its extremities. The variation of the loss as a function of the wavelength for four different waveguide widths has been plotted in Fig. 2. Overall, we see excellent agreement between measurement and simulation. We did observe a discrepancy between the designed waveguide width ($W_D$) and the actual waveguide width ($W$). The waveguides are designed to be 680 nm, 700 nm and 720 nm wide but are found to actually be 710 nm, 730 nm and 755 nm wide respectively. This is confirmed by SEM measurements.

3. Liquid crystal cladding

3.1. Device fabrication

We now turn our attention to LC clad waveguides. In order to facilitate the deposition of a uniform LC layer on the waveguides, a LC cell is assembled based on the technology used in LC display research [13]. The cell consists of an SOI chip (on which the waveguides are lithographically defined) and a glass plate [14]. Since the waveguides are 1 cm long, the glass plate is cut so that it is 7 mm wide. As such it is possible to cover the waveguides and still leave the grating couplers in air for easy coupling of light in and out of the waveguides. The glass plate has a thin layer of indium tin oxide (ITO) deposited on it. The ITO is transparent and acts as the top electrode of the cell while the silicon substrate is the bottom electrode. Accurately controlling the alignment of the LC over the waveguides is very important for the proper operation of our device. The photo-alignment [15] method is an excellent candidate for aligning the LC on top of the waveguides. It possesses many advantages over the rubbing alignment method. Tests were performed with a photo-alignment layer deposited on top of the waveguides. The material was illuminated with UV light. This resulted in LC being aligned parallel to the waveguides. However, measurements revealed that the loss of the waveguides with photo-aligned LC was too high and it was impossible to resolve minima in leakage loss. This is not very surprising given the leaky nature of the TM polarized waveguide modes. Consequently, we resorted to using
the rubbing alignment method. A 2.5 μm thick poly methyl methacrylate (PMMA) layer is spin coated on the ITO. It ensures that the vertical separation between the ITO and the waveguide is large enough to prevent light absorption. In order to preferentially align the director of the LC molecules in a given direction, a thin nylon alignment layer is spin coated onto the PMMA layer. The alignment layer is then rubbed with a soft cloth and forces the LC molecules to align in a planar manner in the direction of the rubbing, with a pretilt of about 2° with the z-axis. We choose the rubbing direction to be parallel to the waveguides on the SOI chip. The glass plate and the SOI chip are glued together with UV curable glue mixed with silica spacer balls. The diameter of the spacer balls determines the thickness of the gap between the glass plate and the chip. Wires are soldered to the ITO on the glass plate and the silicon substrate; this makes it possible to apply a voltage over the LC layer in the cell. The cell is filled with the commercially available LC E7 (ordinary index \( n_o = 1.5024 \) and extraordinary index \( n_e = 1.697 \) at 1550nm [16]) by capillary action. For thin cells, the filling must be performed in vacuum in order to avoid the formation of air bubbles in the cell. A schematic of the fully assembled cell is shown in Fig. 3.

Fig. 3. Schematic of the cell

3.2. Liquid crystal switching behavior

The alignment of the LC is critical to the performance of our device. If the LC molecules covering the waveguides are not well aligned, several domains and dis-inclinations are formed. Light propagating through such a LC layer is scattered at the domain boundaries. Given that lateral leakage loss occurs at the side walls of the waveguides, this scattering will increase the overall loss from the waveguides; making the observation of the lateral leakage loss difficult. Conversely, if the LC molecules covering the waveguides are perfectly aligned, the layer has a well defined director. In this case, the lateral leakage loss should be easier to measure since the scattering by a well aligned LC layer should be negligible.

Figure 4 shows the cell placed in a polarization microscope under crossed polarizers in reflection mode. For zero applied voltage (Fig. 4(a)), we observe that the LC close to the waveguides is well aligned (dark areas) whereas defects form at the edge of the side cladding which are patterned with period structures for pattern density control. The rubbing of the alignment layer on the glass plate is parallel to the waveguides. The LC molecules close to the glass plate align themselves following the rubbing. Close to the waveguides the LC molecules find it energetically favourable to align themselves with their long axes parallel to the waveguides [17].
Fig. 4. LC cell under crossed polarizers in reflection microscopy with lines showing the orientation of the polarizer (P) and analyzer (A). (a) 0Vpp (b) 40Vpp, onset of switching in the LC cell (c) 80Vpp.

results in a uniform alignment in this area of the LC layer. At the side of the etched cladding, the LC molecules are twisted (both left-handed and right-handed) by the periodic structures. Hence we have defects forming over the sides. As the voltage over the LC layer is increased, the defects formed at the sides grow and propagate throughout the cell. Figure 4(b) shows the onset of the LC switching in the cell. For higher applied voltages as in Fig. 4(c) we notice that the defects can propagate towards the waveguides, causing the well aligned part of the LC there to shrink. We can see a different domain boundaries between regions with different alignment in between the waveguides.

The thickness of the LC cladding has a strong effect on the overall loss of the waveguides. Measurements for different LC layer thicknesses reveal that a thin LC layer results in less defects. Since we have a 2.5 μm thick PMMA layer on the top glass plate, we can have a thin LC layer and still avoid absorption by the ITO. We fix the thickness of the LC layer in the cell at 5.6 μm. All results presented further on in this work are for such a cell.

3.3. Tunable leakage loss measurements

We now turn our attention to lateral leakage loss measurements of the LC clad waveguides. The waveguide modes now feel either $n_p$ or $n_r$ in the cladding. This means that the effective indices of the waveguide modes increase compared to the air-clad case. Because of the different orientation of the dominant E-field in the TE and TM modes, the increase in $N_{eff,TM}$ is higher than that in $N_{eff,TE}$. However, $n_{eff,TE}$ is still larger than $N_{eff,TM}$, hence the numerator in (1)
decreases. Accordingly, the wavelength at which the minimum in leakage loss occurs (magic wavelength) shifts to shorter wavelengths as depicted in Fig. 5.

We note from Fig. 5 that the leakage loss exhibited in air-clad waveguides is preserved in the LC clad waveguides. The waveguide with \( W_D = 740\,\text{nm} \) is interesting since it exhibits its magic wavelength within the wavelength window of interest. The leakage loss variation of the other waveguides (with different width) indicate that their magic wavelength lies at shorter \( (W_D = 720\,\text{nm}) \) and longer \( (W_D = 780\,\text{nm}) \) wavelengths respectively. We have modeled these LC clad waveguides and found good agreement between our simulations and the measurements as depicted in Fig. 6(a). Note that the theoretical loss minimum is 0dB/cm which is obviously not the case in the practice. The waveguides designed to be 720nm and 740nm wide are found to correspond to 750nm and 785nm wide waveguides respectively. In order to determine the range over which the magic wavelength can be tuned, we model the 5.6\,\mu m thick cladding layer as a uniaxial material with c-axis along the z (green dashed curve) and y (black dashed curve) axes respectively. With the former and latter being the initial and final states of our LC cladding. Figure 6(b) shows a plot of the E-field components in the waveguide. Notice that in the area occupied by the LC in the cell (y distance greater than 0.11\,\mu m), the y component is stronger than the z component. Accordingly, the position of the magic wavelength can be switched by more than 70nm by reorienting the LCs in the cladding from being aligned along the z axis to being aligned along the y-axis (see Fig. 6(a)).

The tuning of the leakage loss behavior is achieved by applying a voltage over the LC cladding. This causes the LC molecules to reorient themselves so that they are more and more parallel with the electric field lines (along the y-axis). The index felt by the TM-like mode in the upper cladding is given by:

\[
n_{clad} = \frac{1}{\sqrt{\left(\frac{\cos \alpha}{n_e}\right)^2 + \left(\frac{\sin \alpha}{n_e}\right)^2}} \tag{2}
\]

with \( \alpha \) the angle the LC director makes with the z-axis. For zero applied voltage, \( \alpha = 0 \).
increases with increasing applied voltage attaining its maximum value of 90° for high applied voltages. Accordingly, \( n_{\text{clad}} \sim n_o \) for low voltages and \( n_{\text{clad}} \sim n_e \) for high voltages. The applied voltage is swept from 0Vpp to 300Vpp in steps of 20Vpp. For each voltage step, we measure the leakage loss over the wavelengths ranging from 1510nm to 1590nm. We plot the result of these measurements for the 785nm wide (\( W_D = 740\)nm) waveguide in Fig. 7(a). As the voltage is increased, the wavelength at which the minimum in leakage loss occurs shifts to shorter wavelengths. We measure a shift from 1564nm at 0Vpp to 1524.5nm at 300Vpp; this is about 56 % of the change predicted from simulations for perfectly aligned LC cladding layers. We also notice that the shift in the wavelength at which the minimum in leakage loss occurs is accompanied by an increase in loss. On average the loss in the waveguide increases by about 8dB/cm when the voltage is changed from 0Vpp to 300Vpp.

![Fig. 6. (a) Simulated versus measured loss values for LC clad waveguides. Solid lines: measured lateral leakage loss data. Dashed lines: simulation data. Dark dashed line: y-aligned cladding. (b) Field profiles of the various E-field components for a \( W = 785\)nm waveguide with air cladding. Notice that the y component is much stronger than the z component in the area (y distance greater than 0.11 \( \mu m \)) occupied by the LC.](image)

![Fig. 7. (a) Voltage tuning of the loss in a \( W_D = 740\)nm wide LC clad waveguide. The voltage is ramped from 0Vpp (black) to 300Vpp (yellow) in steps of 20Vpp. The arrow indicates the direction in which the loss curve shifts when the applied voltage is increased. (b) Voltage dependence of the lateral leakage loss in a \( W_D = 740\)nm LC clad waveguide.](image)
We complete this section by taking a look at the voltage dependence of the leakage loss of the LC clad waveguide. Figure 7(b) shows plots of this dependency for five different wavelengths. For 1550nm, 1570nm, and 1580nm, the leakage loss increases with increasing voltage. For example at 1570nm the loss increases from 4.5dB/cm at 0Vpp to 18.1dB/cm at 300Vpp. Whereas at 1580nm the loss increases from 5.7dB/cm at 0Vpp to 21.5dB/cm at 300Vpp. For 1520nm and 1530nm, the loss decreases with increasing voltage but as the voltage increases past a certain point, it increases again. For example at a wavelength of 1520nm, the leakage loss decreases from 12.2dB/cm at 0Vpp to 9.2dB/cm at 140Vpp. For voltages higher than 140Vpp, the loss starts increasing again.

4. Discussion

The measurements for the LC clad waveguides are only corrected for the contribution of the input and output grating couplers. We do not take into account the difference in loss between the parts of the waveguide covered (70%) and uncovered (30%) by the glass. Observation of the cell under the polarization microscope reveals that LC also covers the part of the waveguides uncovered by the glass plate. As a result we have the part of the waveguides uncovered by glass covered with LC oriented along the z-axis. The LC over this part of the waveguides does not reorient with increasing applied voltage whereas the LC in the glass covered part does. Accordingly, the loss we measure is a weighted average of the contribution from both parts of the waveguide according to the formula;

\[
Loss = 0.3Loss_{LC, uncovered} + 0.7Loss_{LC, covered}
\]

In order to verify this, we model the \( W = 785\)nm waveguide with an increasingly reoriented LC upper cladding layer. This is achieved by using \( n_{clad} \) for the cladding layer above the waveguide and varying the value of \( \alpha \) from 0° to 70° in steps of 10°. The case \( \alpha = 0 \) corresponds to the measurement at 0Vpp. When we apply (3) to the simulation results we obtain a plot similar to Fig. 7(a).

![Fig. 8. Lateral leakage loss as a function of wavelength for increasing \( \alpha \).]
Figure 8 exhibits the same features as Fig. 7(a); as $\alpha$ (i.e. the applied voltage) increases, the wavelength at which the minimum in lateral leakage loss occurs experiences a blue shift. The wavelength shift is accompanied by an increase in the lateral leakage loss. This only qualitatively explains the measurements for the LC clad waveguides. Closer inspection reveals that this effect alone cannot quantitatively explain the increase in loss. The LC molecules respond to an applied voltage by not only tilting (in the yz plane), but also twisting (in the xz plane). This results in fluctuations in the LC director profile along the length of the waveguide. Each LC director profile corresponds to a given phase-matching angle between the radiating slab TE mode and the guided TM-like mode. Since each phase-matching angle corresponds to a different magic wavelength, the measured loss is a weighted average of the loss corresponding to each of the different LC director profiles which occur along the length of the waveguide. The overall result is an increase in measured loss with increasing applied voltage.

The performance of the device can be improved by minimizing the effect of the defects generated by the sides. For low applied voltages, this can be achieved by increasing the separation between the waveguide and the sides. When a voltage is applied to the LC layer, the defects propagate through the cell. Given that defects in a LC cell typically extend over a distance comparable to the thickness of the cell, the scattering in the LC layer close to the waveguide increases with increasing applied voltage. Increasing the distance between the sides and the waveguide will also make it possible to achieve a greater reduction in the scattering with increasing voltage than achieved here.

The voltages mentioned above are quite high. The LC cells we fabricate have several dielectric layers between the two electrodes (see Fig. 3). The dielectric permittivity of silica, PMMA and LC in the KHz range is equal to 3.9, 2.6 and 5.1 (for low voltages, and 19.6 for high voltages). Accordingly, despite the fact that the LC layer is thicker than the other dielectric layers, the voltage drop over it is only a fraction of the total applied voltage. For low and high applied voltages this fraction is equal to about 40% and 15% of the applied voltage respectively. The problem of high applied voltages can be solved by doping the shallow-etched waveguide slab, so that it can be used as an electrode.

The method we propose for tuning the position of the leakage loss minima of LC clad shallow-etched waveguides opens up the possibility to have voltage-tunable reconfigurable single-mode optical interconnects (waveguides) operating over a wide band of wavelengths. As we show, such interconnects can be wider than other silicon photonic single-mode waveguides and hence potentially more fabrication tolerant. We have also demonstrated wavelength selective transmission with an extinction ratio of 20dB/cm. More sophisticated waveguide designs can be conceived that would increase the number of wavelengths which are efficiently transmitted through a single waveguide, and thus single-mode on-chip wavelength division multiplexing should be possible. Our work also opens up interesting possibilities for applications in which it is important to tune the loss of a signal. A high extinction Mach-Zehnder interferometer is an example of such an application. If you can balance the power exactly between the arms by trimming the loss (using the LC) in one arm, you can get higher contrast interference. It should also be noted that the power lost from the waveguide through the lateral leakage phenomenon is converted into coherent TE radiation. By controlling the rate of leakage along the waveguide, it is possible to achieve on-chip beam forming [18]. Our demonstration of dynamic reconfiguration of lateral leakage using LCs opens up the possibility for dynamic beam forming on-chip with applications in information imaging and sensing as well as information processing and signal routing.
5. Conclusion

We have demonstrated a new method for tuning the lateral leakage loss of TM-like modes in shallow etched LC clad SOI waveguides. We started out by measuring the leakage loss in an air-clad waveguide. We then proceed by giving details about how to incorporate a LC cladding on the waveguide. We then measure the leakage loss of the LC clad waveguides and find it to be comparable to the air-clad case. Finally, we also show that for a fixed wavelength, the leakage loss can be modulated by appropriately modulating the applied voltage.

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Chapter 3

Long-range photonic bus

3.1 Long-range coupling of silicon photonic waveguides using lateral leakage and adiabatic passage

Long-range coupling of silicon photonic waveguides using lateral leakage and adiabatic passage

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Abstract: We present a new approach to long range coupling based on a combination of adiabatic passage and lateral leakage in thin shallow ridge waveguides on a silicon photonic platform. The approach enables transport of light between two isolated waveguides through a mode of the silicon slab that acts as an optical bus. Due to the nature of the adiabatic protocol, the bus mode has minimal population and the transport is highly robust. We prove the concept and examine the robustness of this approach using rigorous modelling. We further demonstrate the utility of the approach by coupling power between two waveguides whilst bypassing an intermediate waveguide. This concept could form the basis of a new interconnect technology for silicon integrated photonic chips.

References and links
1. Introduction

Mass manufacture of monolithic systems of extraordinary complexity, compactness and precision using CMOS processing has underpinned the information revolution. However, interconnections between complex functional blocks remains a critical challenge, often requiring numerous interconnect layers above the functional plane. The CMOS process has recently been adapted to photonic integrated circuits [1] with applications emerging in high-speed communications [2], photonic signal processing [3] and quantum optics [4]. Silicon photonic systems are gaining momentum but device complexity will again be limited by interconnect technology. Out of plane optical interconnect techniques have been proposed [5], but these are not compatible with emerging CMOS silicon photonics standards which permit only a single silicon optical wave guiding layer [1]. In-plane crossing structures, which are CMOS compatible, have been demonstrated [6], but these can introduce losses and may be sensitive to fabrication variations.

Long range communications between waveguides through unguided radiation in the silicon slab have been proposed as an alternate interconnect solution [7]. This particular approach uses thin, shallow ridge silicon on insulator waveguides which, when operated in the TM mode, can radiate into the TE modes of the slab [8, 9]. This TE radiation is traditionally considered a loss mechanism, however as it is a coherent process, with appropriate control over the radiation it could be utilised as a resource [10]. We have previously shown that it is possible to control the radiation direction [7] and also generate directed, collimated beams [11]. The nature of this radiation is quite sensitive to the waveguide geometry and thus may not be robust to fabrication variations. Further, if this radiation is to be used as an interconnect, then the unbound nature of this radiation may lead to undesired interaction with intermediate functional blocks.

Adiabatic techniques are well known in photonics, principally being invoked when properties of a single waveguide or two waveguide system are changed slowly, for example with adiabatic tapers [12]. Slow changes can also be used to effect population transfer between waveguides through a technique called Coherent Tunnelling Adiabatic Passage (CTAP) which is a spatial analogue of the well-known STIRAP (STImulated Raman Adiabatic Passage) protocol in quantum optics [13]. CTAP was originally proposed for massive particles in tight-binding systems [14, 15] and then extended to waveguides [16, 17]. Transfer via multiple intermediate states has also been considered [18–21]. CTAP has the advantage that the transport is extremely robust against fluctuations in the coupling between sites. CTAP also has the surprising feature that the population in the intervening site is greatly suppressed, and in the adiabatic, tight-binding limit, is identically zero. This unusual behaviour raises the question of whether CTAP may be exploited to achieve robust long range coupling between waveguides via unbound radiation, but without exciting this radiation.

Here we propose and numerically demonstrate the combination of CTAP and lateral leakage to achieve a new type of coupler. Light guided within one waveguide can be transferred over a long distance to another waveguide through use of an unbound lateral leakage state which is coupled to both waveguides. Due to the nature of CTAP, this coupling is extremely robust, being relatively independent of coupling length and remarkably, the intermediate radiation is not populated during the coupling. We also show that this technique can be used to bypass an intermediate waveguide without cross-talk.

This paper is organised as follows: Section 2 presents a brief overview of the CTAP protocol in the context of optical modes and Section 3 reviews lateral leakage and shows how coupling between bound waveguides modes and lateral leakage radiation can be controlled. Section 4 numerically simulates the CTAP coupling between waveguides and tests the robustness of this technique with varying device length. Section 5 then shows the bypass of an intermediate waveguide. Finally, Section 6 discusses the limitations of this specific demonstration and outlines the opportunities for future research on this approach.
2. Coherent tunnelling adiabatic passage

Coherent Tunnelling Adiabatic Passage (CTAP) is a protocol for transferring population between defined states. In particular, the transport should be spatial. It is usual that the modes be in some sense equivalent or discrete, however such restrictions are not always necessary.

To illustrate CTAP, consider a three-state system as shown in Fig. 1(a). The states $|L\rangle$ and $|R\rangle$ are mutually isolated and can only couple to the common state $|B\rangle$, which acts as a bus. The strength of the couplings between each state and the bus are $\Omega_L$ and $\Omega_R$.

The Hamiltonian describing this problem is

$$H(z) = \sum_{i=L,B,R} \beta_i |i\rangle\langle i| + \Omega_L |L\rangle\langle L| + \Omega_R |R\rangle\langle R| + h.c.,$$

where $\beta_i = k_0 n_i$ is the propagation constant for mode $i$ with effective index $n_i$, and $k_0$ is the propagation constant of the free space.

The CTAP protocol is achieved when the couplings are varied in the so called counter-intuitive sequence. This requires $\Omega_L(0) \gg \Omega_R(0)$, and gradual variation in each with increasing $z$ until $\Omega_R(z_{max}) \gg \Omega_L(z_{max})$. There is considerable flexibility in the actual sequence implemented, and popular choices include Gaussian [22] and sinusoidal [23] variations, although discontinuities in the controls can also be tolerated under certain conditions [24, 25]. Here we choose squared sinusoid as per Fig. 1(b). The counter-intuitive sequence works by maintaining the system in the null state, which is the supermode (in the limit that all of the $\beta_i$ are equal):

$$|D_0\rangle = \frac{\Omega_R |L\rangle - \Omega_L |R\rangle}{\sqrt{\Omega_L^2 + \Omega_R^2}}.$$

Note that this has the desired properties for adiabatic passage, namely that when $\Omega_L \gg \Omega_R$, $|D_0\rangle = |R\rangle$, and when $\Omega_R \gg \Omega_L$, $|D_0\rangle = |L\rangle$. Provided adiabaticity is preserved, the population in $|B\rangle$ will be identically zero, although the population in $|B\rangle$ only approaches zero when finite mode size is taken into account [26]. Here, adiabaticity is defined with respect to the separation (in terms of energy) between $|D_0\rangle$ and the nearest supermode. Hence the scheme is largely immune to small errors in realisation. It is also important to recognise that the system is highly insensitive to loss or decoherence mechanisms that act on the bus state, due to the suppressed population there [27–29].

3. Thin shallow ridge waveguides and control of lateral leakage

Having introduced CTAP in Section 2, this section introduces thin shallow ridge waveguides and lateral leakage behaviour and shows how this leakage can be controlled for the purpose of implementing a CTAP coupler with this system.

Fig. 1. The CTAP protocol taking population from $|R\rangle$ to $|L\rangle$: (a) A 3 state scheme with two isolated states coupled to a central bus, (b) counter-intuitive evolution of coupling strengths $\Omega_L$ and $\Omega_R$, (c) population evolution in states $|L\rangle$, $|R\rangle$ and $|B\rangle$. 

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3.1. Lateral leakage from thin shallow ridge waveguides

Thin shallow ridge waveguide can be realised using standard CMOS processing and offer highly evanescent modes with low propagation losses. The TM mode can exhibit leakage of power into the laterally radiating TE slab mode. However, this leakage occurs only at the waveguide side walls and at so-called ‘magic’ waveguide widths, the radiation from each side wall cancels [8].

Consider the thin shallow ridge structure of Fig. 2(a). To examine the isolated modes of this structure, the waveguide widths were set to the magic width of 0.7 μm at a wavelength of 1550 nm. The guided modes of the system were simulated using a mode matching method [30]. Three simulated guided modes of the system are presented in Fig. 2(b)-(d). Each had the same effective index and were thus degenerate. Figure 2(b) and (c) present the isolated TM modes |L⟩ and |R⟩ respectively. Figure 2(d) presents the TE slab ‘bus’ mode |B⟩.

For simplicity, the slab has been terminated, as illustrated in Fig. 2(a), and thus the TE slab radiation is in fact a discrete mode with an oscillating standing wave pattern.

To illustrate the impact of coupling, the widths (W_R) of waveguides |L⟩ and |R⟩ were set to 1.22 μm, such that they were strongly and equally coupled to the TE slab |B⟩. The modes of the system were again simulated using mode matching and the resulting supermodes, corresponding to the eigenstates of Eq. (1), are presented in Figs. 2(e)-(g). Fig. 2(f) is a supermode with equal population in each of the TM modes |L⟩ and |R⟩ and no population in |B⟩, i.e. the null state, |L⟩ − |R⟩. Whilst Figs. 2(e) and (g) are the supermodes |L⟩ ± √2|B⟩ + |R⟩ with strong population in the TE slab |B⟩. The three modes of Fig. 2(e)-(g) are no longer degenerate as the coupling has caused significant splitting of the effective indices of the three modes.

As discussed in Section 2, CTAP requires adiabatic transformation of the coupling to transfer population from |R⟩ of Fig. 2(c) at the start, into the coupled supermode of Fig. 2(f) in the middle, and then into |L⟩ of Fig. 2(b) at the end. One might consider simply tapering the width of the waveguides to control the coupling, as demonstrated in [11], however, the modal effective index is sensitive to the waveguide width. For optimal CTAP, it is important that the effective indexes of |L⟩ and |R⟩ remain equal. Hence an alternative coupling approach is required.
3.2. Control of lateral leakage using waveguide location

An approach to controlling the coupling between the TM waveguides and TE slab that will maintain equal effective indexes for the modes $|L\rangle$ and $|R\rangle$ is suggested by the standing wave pattern of the TE mode as illustrated in Fig. 2(d). It might be expected that the coupling between the guided TM mode and the TE slab should depend strongly on the lateral location of the thin shallow ridge waveguide. To establish the effect of waveguide location on coupling between the TM mode and TE slab, the structure of Fig. 3(a) was modelled. A single thin shallow ridge waveguide was located on a broad slab. The width of the thin shallow rib was set to 1.22 $\mu$m such that the TM guided mode should, in principle, be strongly coupled to the TE radiation.

For the particular CTAP protocol we aim to implement it is necessary that $|L\rangle$, $|R\rangle$ and $|B\rangle$ all have the same effective index and are hence degenerate when uncoupled. Referring to Fig. 3(a), the waveguide was placed in the centre of the slab, and the slab width was adjusted to find a configuration where the the TM mode and TE slab are degenerate and uncoupled. Mode matching was used to simulate the effective index of the TE and TM modes of this structure as a function of slab width. The results are presented in Fig. 3(b). The TE slab mode effective indexes vary with slab width while the index of the TM guided mode remains almost constant. When the TM and TE modes are degenerate, if the symmetry is not matched, the indexes simply cross; however, if the symmetry matches, mode splitting occurs leading to an anti-crossing. Figure 3(b) shows that it is possible to select a slab width where there is a TM guided mode and TE slab mode that are degenerate, but uncoupled at a slab width of 30.4 $\mu$m.

Next the impact of waveguide location on coupling between the TE and TM modes was investigated. The location of the thin shallow ridge of Fig. 3(a) was translated laterally across the slab and mode matching was used to simulate the effective indexes of the two supermodes of the system as a function of waveguide offset. The results are presented in Fig. 3(c). At 0 nm displacement, the modes are degenerate and uncoupled. As the waveguide was translated, the indexes split, indicating coupling, reaching a maximum at a displacement of 370 nm. Further displacement decreased the mode splitting until degeneracy was again reached at 740 nm corresponding to a half cycle of the standing wave pattern allowing us to define a coupling period $\zeta = 740$ nm. These results show that it is indeed possible to control the coupling between the TM and TE modes using waveguide location and this technique could be utilised to implement CTAP with these waveguides.

4. Demonstration of long range coupling using CTAP and lateral leakage

Section 3.2 established that it is possible to control the coupling between localised waveguide modes and distributed slab modes by adjusting their locations. We now show how this coupling
control technique can be used to implement a CTAP protocol with thin shallow ridge waveguides. Specifically, it is shown that power can be adiabatically transferred between two isolated waveguides using TE slab mode radiation as an intermediate bus, but without ever populating this bus. This section will also test the robustness of this approach by exploring the impact of adjusting device length on the propagation.

4.1. CTAP using Lateral Leakage

Optical propagation in the longitudinally varying structures of this section were simulated using eigenmode expansion (EME) [7]. EME was chosen as the most appropriate model as Beam Propagation Method (BPM) cannot handle large TE radiation angles and Finite Difference methods have expensive computational requirements for long devices. The EME model rigorously treats the waveguide translation and its impact on the underlying supermodes of the system accounting for radiation. It would be beneficial to validate these findings using a software tool that does not assume modes, such as Finite Difference Time Domain (FDTD). However such solutions are difficult in the case of thin-ridge devices due to the necessity to keep track of features at many different length scales, including the nanometer scale of the thin-ridge of the waveguide, the micron scale of the light, waveguide width and separation, and the millimeter scale of the total device.

Figure 4(a) presents the cross-section of the geometry under consideration. Two thin shallow ridge waveguides supporting TM modes $|L\rangle$ and $|R\rangle$ were placed on a silicon slab supporting a distributed TE slab mode $|B\rangle$. The waveguide widths ($W_R$) were 1.22 $\mu$m such that the TM and TE modes should be coupled as shown in Section 3.2. The slab width was adjusted from 30.4 $\mu$m to 30.78 $\mu$m to account for an additional waveguide while ensuring $|L\rangle$, $|R\rangle$ and $|B\rangle$ are degenerate when isolated. The location of the two waveguides were adjusted to control the coupling between the modes. Light was coupled into and out of the system through short sections of non-radiating magic width waveguide of width 0.70 $\mu$m.

The approach for determining the strategy for translating the pair of waveguides is as follows. Although the system has been altered by the presence of an additional waveguide, since the modes of the system have been tuned to match that of Section 3.2, it would be expected that the
Fig. 5. (a) plan view of longitudinally invariant waveguides (pink indicates strong coupling, green indicates weak coupling), excitation on $|R\rangle$; (b) optical propagation for uncoupled configuration; (c) plan view of translated waveguides in intuitive CTAP configuration, excitation on $|L\rangle$; (d) optical propagation for intuitive configuration; (e) plan view of translated waveguides in counter-intuitive CTAP configuration; (f) optical propagation for counter-intuitive configuration. (g) supermode effective index throughout CTAP evolution; In each case $z_{\text{max}} = 10\text{mm}$.

coupling behaviour of the modes would oscillate in a similar fashion with waveguide translation as observed in Fig. 3(c). Both waveguides were translated symmetrically away from the centre of the slab and the effective index of the three modes close to the TM guided mode of an isolated waveguide were found using the same mode-matching eigensolver as used in Section 3.2. Figure 4(b) presents the effective index of each mode as a function of waveguide separation and oscillatory behaviour is clearly evident with a period of $\zeta = 740$ nm in separation, similar to Fig. 3(c).

The maximum splitting, corresponding to strongest coupling between the two waveguides $|L\rangle$ and $|R\rangle$ and the slab mode $|B\rangle$, is found to be at a central core separation of 6.6 $\mu$m. Unlike Fig. 3(c), there are now 3 modes in this system and it would appear that one mode remains unperturbed throughout the transition. It is expected that this mode corresponds to strong coupling between the two TM modes of $|L\rangle$ and $|R\rangle$ waveguides but complete isolation from the slab mode $|B\rangle$. This is the mode that we would want to populate during the CTAP transition. We now wish to find the required relative offsets for the start and end of the transition. These would be characterised by complete independent isolation of $|L\rangle$ and $|R\rangle$.

To characterise the impact of relative waveguide offset and coupling of the waveguides to each other, the two waveguides were placed symmetrically at the maximally coupled separation of 6.6 $\mu$m, the left waveguide $|L\rangle$ was held stationary and the right waveguide $|R\rangle$ was further translated and the proportion of TM field in both $|L\rangle$ and $|R\rangle$ waveguides was assessed for the central, unperturbed mode. Figure 4(c) presents the square of the magnitude of the vertical electric field component $|E_y|^2$ for this mode in the regions of each waveguides $|L\rangle$ and $|R\rangle$ as a function of the location of waveguide $|R\rangle$. The field $|E_y|^2$ is taken as a measure of the presence of the TM mode in each waveguide. It can be seen that at the maximally coupled state, with both $|L\rangle$ and $|R\rangle$ at 3.3 $\mu$m, the TM field is evenly distributed between the two waveguides as was predicted. As $|R\rangle$ is translated, the TM field in $|L\rangle$ gradually decreases reaching a null when $|R\rangle$ is located at 3.47 $\mu$m. Further translation of $|R\rangle$ shows the sequence repeating with a period
of $\zeta/2 = 370$ nm as expected. Thus when $|L\rangle$ is at 3.3 $\mu$m and $|R\rangle$ and $|B\rangle$ will be strongly coupled and $|R\rangle$ will be isolated.

The configuration of Fig. 5(a) was considered first with $|L\rangle$ located at $x=-3.30$ $\mu$m to be strongly coupled to $|B\rangle$; and $|R\rangle$ located at $x=+3.47$ $\mu$m such that it is isolated from $|B\rangle$ throughout propagation. The separation between waveguides of 6.77 $\mu$m will be sufficient to ensure no appreciable evanescent coupling directly between $|L\rangle$ and $|R\rangle$.

A simulation was performed with $|R\rangle$ excited as indicated by the red arrow on Fig. 5(a). The simulation results are presented in Fig 5(b) showing excitation of the fundamental TM mode of the isolated waveguide with minimal radiation loss from the input region to the propagation region, and no evidence of coupling to either the TE slab $|B\rangle$ or the other TM mode $|L\rangle$.

The geometries of Fig. 5(c) and (e) are both identical and designed to provide an acceptable CTAP coupling scheme, with the system response strongly dependent on the initial excitation. At the input $|L\rangle$ was located at $x=-3.30$ $\mu$m (coupled to $|B\rangle$) and $|R\rangle$ was at $x=+3.47$ $\mu$m (isolated from $|B\rangle$). However, during propagation, the locations of $|L\rangle$ and $|R\rangle$ were linearly translated, such that at the output, $|L\rangle$ was offset by -3.47 $\mu$m, (isolated from $|B\rangle$), and $|R\rangle$ was at +3.30 $\mu$m (coupled to $|B\rangle$). From Section 3.2, linear translation corresponds to sinusoidal evolution of the coupling strength. A simulation was performed with $|L\rangle$ excited as indicated by the red arrow on Fig. 5(c). The results are presented in Fig 5(d). At the input, light rapidly couples back and forth between $|L\rangle$ and $|B\rangle$. Mid-way, there is equal and in-phase excitation in both $|L\rangle$ and $|R\rangle$ and rapid coupling to $|B\rangle$ continues with the same coupling length. At the output, the excitation has transferred to $|R\rangle$ with rapid coupling to $|B\rangle$ still evident. The output power is split between $|R\rangle$ and $|B\rangle$. This split will be highly sensitive to device length and has been seen in such systems before [15, 32].

The structure and excitation of Fig. 5(e) should achieve counter-intuitive CTAP coupling. The device geometries are identical to Fig. 5(c), however excitation has been change to $|R\rangle$ as indicated by the red arrow in 5(e). The results are presented in 5(f) which shows smooth transition of the optical power from $|R\rangle$ to $|L\rangle$ without appreciable excitation of $|B\rangle$. Some slight oscillation is evident, however the rapid, oscillatory coupling to $|B\rangle$ seen in 5(e) are not present. The absence of these oscillations is a major distinguishing feature between CTAP and devices such as directional couplers.

Observing the effective indices of the system supermodes throughout CTAP evolution are shown in Fig. 5(g). The counter-intuitive sequence only excites a single mode (the isolated mode with effective index in green in Fig. 5(g)) whereas the intuitive case excites a superposition of two modes (the two coupled modes with effective indexes indicated in blue/red in Fig. 5(g)) which explains the modal beating observed.

4.2. Suppression of bus mode excitation and adiabaticity of long range coupling

To more closely examine the excitation of $|B\rangle$ during the adiabatic transfer from $|R\rangle$ to $|L\rangle$, the $E_x$ component of the results of Fig. 5(f) were replotted corresponding to the TE polarisation. These results are presented in Fig. 6(a). It is evident that there is, in fact, some slight excitation of $|B\rangle$. There are several effects that can contribute to this residual excitation, including the staircase approximation [25], finite spatial extent of the modes [26], residual non-adiabaticities in the evolution [15], imperfect initialisation in the null state [21] and imperfect coupling of power to the modes of the system at the input and output of the structure. Figure. 6(a) suggests that imperfect coupling to $|R\rangle$ at the input is the dominant source of the population in the TE mode $|B\rangle$, but this effect is deemed negligible for the current demonstration.

The robustness of CTAP protocol was explored by monitoring the coupled power while varying the total device length. Once in the adiabatic regime, the transport was expected to be largely independent of the exact device length, asymptotically approaching perfect transport. This con-
Fig. 6. (a) Lateral ($E_x$) optical propagation of device above adiabatic limit ($z_{max} > A_{lim}$); (b) $|L\rangle$ output power as a function of device length ($z_{max}$); (c) Optical propagation of device with $z_{max} < A_{lim}$; (d) TE polarised field ($E_y$) close to the input termination (common to all simulations).

Contrasts non-adiabatic couplers where the final power would depend critically and periodically on the device length relative to the coupling length. The structure of Fig. 5(e) was simulated, but with device length varied from $z_{max} = 0.5$ to 10 mm in steps of 50 $\mu$m. Figure 6(b) presents the power coupled from $|L\rangle$ at the output as a function of device length. For lengths of 2 to 10 mm, the output remains relatively constant indicating adiabatic behaviour while we are operating above the adiabatic limit ($z_{max} > A_{lim}$) where for this particular structure, $A_{lim}$ is around 1 mm as indicated in Fig. 6(b). The transmission is slightly less than unity and there is a slight ripple evident in the transmission as the length is varied which could be due to the imperfect coupling mentioned above. When the length drops below $A_{lim}$, the transmission begins to drop, falling off dramatically for lengths below 1 mm. This drop off is due to the device being too short to exhibit adiabatic passage.

Figure 6(c) presents the propagation for the structure of Fig. 5(e) with $z_{max} = 500$ $\mu$m. Light input to $|R\rangle$ initially remains isolated from $|B\rangle$, but unlike the behaviour of 5(f), mid-way the light remains in $|R\rangle$ and simply radiates into $|B\rangle$ with minimal coupling to $|L\rangle$. Figure 6(d) presents a highly magnified view of $E_x$ close to the input showing energy naturally radiating from $|R\rangle$ into $|B\rangle$. This coupling occurs at the input where $|R\rangle$ should be isolated, providing evidence that the excitation and isolation of $|R\rangle$ is not perfect.

5. CTAP using Lateral Leakage to bypass an intermediate waveguide

Whilst the demonstration of Section 4 is interesting, this does not provide the functionality for long range interconnections across a complex planar system. We now show this functionality by demonstrating that CTAP using lateral leakage can bypass an intermediate waveguide.

The structure of Fig. 7(a) is similar to that of Fig. 5(a), but has an additional intermediate waveguide, $|I\rangle$, inserted at the centre. The waveguide supporting $|I\rangle$ was maintained at the magic width throughout propagation in order to isolate it from the slab mode, $|B\rangle$ irrespective of its location. The width of the slab was altered to 30.385 $\mu$m to ensure that $|B\rangle$ was phase matched to $|L\rangle$ and $|R\rangle$. The offset on $|L\rangle$ and $|R\rangle$ were $\pm 5.652$ $\mu$m to achieve isolation and $\pm 5.484$ $\mu$m to achieve coupling to $|B\rangle$. This increased offset from the centre aimed to ensure no evanescent coupling between $|L\rangle$, $|R\rangle$ and $|I\rangle$. The structure was configured as in Fig. 7(b) such that at the input, $|R\rangle$ was isolated and $|L\rangle$ was coupled to $|B\rangle$ and followed the same counter-
intuitive translation as in Fig. 5(f). The intermediate state $|I\rangle$ was not translated, however it would be expected that translation of the intermediate waveguide would not impact the performance of the device. Each of the three waveguides was interfaced at the input and output to non-radiating magic width waveguides. The structure was simulated as described in Section 4.

The first simulation tested the isolation of $|I\rangle$ from $|L\rangle$, $|R\rangle$ and $|B\rangle$. Optical power was input to the intermediate waveguide as indicated by the red arrow in Fig. 7(b). Figure 7(c) presents the simulated results showing that light remains confined to the intermediate waveguide without any evidence of coupling. Next the intuitive coupling case of Fig. 7(d) was simulated and is presented in Fig. 7(e). These results can be compared to Fig. 5(e) exhibiting similar population oscillations. Importantly, there is no evident coupling into $|I\rangle$, as expected since it is at the magic width and should be isolated from the TE slab.

Finally, the counter-intuitive coupling case of Fig. 7(f) was simulated and the results are presented in Fig. 7(g). Comparing these results to Fig. 5(f), it can be seen that again adiabatic passage without appreciable population in either the bus mode, $|B\rangle$, or intermediate waveguide, $|I\rangle$, has been achieved. Slight pulsing of the light is again observed in Fig. 7(g) similar to that of Fig. 5(f). These simulations confirm that this adiabatic coupling structure is indeed capable of transferring an optical signal from one waveguide to another, bypassing an intermediate waveguide using the TE slab mode as a type of bus, but without ever populating this bus.

6. Conclusions

We have described a new concept for adiabatic transfer of power between two thin shallow ridge waveguides and proved this concept using rigorous numerical simulation. The power transfer occurs by coupling each waveguide to a laterally distributed slab mode which acts as an optical bus. The novelty of our demonstrated concept is that due to the nature of the Coherent Tunnelling Adiabatic Passage (CTAP) protocol employed, power is robustly transferred from one waveguide to the other without ever populating the intermediate optical bus. The distributed nature of the bus allows the coupling to be long-range, exceeding evanescent interaction distances and indeed extending beyond nearest neighbour interactions. We have demonstrated...
this feature by showing that our CTAP coupler can act as a cross-connect bypassing an intermediate waveguide which is immune to cross-talk. Since the bus state population is minimal, and the CTAP protocol is highly robust we would expect the transfer to be insensitive to other intervening structures or imperfections of the slab.

We propose that this new coupling technique could have a potential application as an interconnect mechanism across complex integrated optical systems. However, our initial device can benefit from further optimization before this approach can be taken beyond the proof of concept stage, there are limitations and possible extensions that should be explored.

A significant restriction of our demonstration is that in our demonstration the transport was via a discrete mode of the slab. This has obvious limitations as it imposes a restriction on the properties of the whole slab, rather than just the slab in the vicinity of the active waveguides. However, there are STIRAP/CTAP protocols that operate using multiple intermediate states [37, 38] and even via a continuum [39], again with minimal occupation of those intermediate states. Since our CTAP protocol is a direct analogy of STIRAP we are confident that similar approaches could be employed to eliminate the dependence on the properties of the discrete modes of the slab. Fabrication tolerances can be improved by further waveguide engineering [31,40] and can provide enhanced coupling to reduce overall device lengths. Before fabrication it would be beneficial to validate these findings using a software tool that does not assume modes, such as finite difference time domain, but this is extremely challenging and is thus proposed as future work.

In our demonstration of adiabatic transfer bypassing an intermediate waveguide only two waveguides were coupled to the slab at any one time with the third intermediate waveguide maintained at the magic width and hence uncoupled from the slab at all times. It would be of interest to explore cases where more than two waveguides are coupled to the bus simultaneously, for example topologies equivalent to the tripod and multi-pod schemes from STIRAP. These schemes have been proposed for applications such as geometric gates [40] and multiple-recipient adiabatic passage [41], which cannot be realised without some form of non-nearest neighbour coupling, such as has been outlined here.

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Chapter 4

Long-range adiabatic quantum gates suited for integrated photonics

4.1 Adiabatic two-photon quantum gate operations using a long-range photonic bus

Adiabatic two-photon quantum gate operations using a long-range photonic bus

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Abstract

Adiabatic techniques have much potential to realize practical and robust optical waveguide devices. Traditionally, photonic elements are limited to coupling schemes that rely on proximity to nearest neighbour elements. We combine adiabatic passage with a continuum based long-range optical bus to break free from such topological restraints and thereby outline a new approach to photonic quantum gate design. We explicitly show designs for adiabatic quantum gates that produce a Hadamard, 50:50 and 1/3:2/3 beam splitter, and non-deterministic controlled NOT gate based on planar thin, shallow ridge waveguides. Our calculations are performed under conditions of one and two-photon inputs.

Keywords: quantum computation, silicon photonics, adiabatic transport, quantum information processing, two photon, single photon, gates

1. Introduction

Integrated optics is becoming one of the most important platforms for the production of compact, scalable, linear optical quantum devices [1]. Much of this progress derives from the use of laser-defined waveguides in glass or polymer, which enables compact three-dimensional waveguide geometries to be designed and rapidly prototyped [2]. While these devices are clearly important for scientific applications, they are not compatible with standard lithographic processes and are limited in the topologies that can be considered. Direct write waveguides rely on serial material modification which places limitations on the complexity of the waveguide system designs that can be achieved. As a result, most geometries rely on evanescent coupling of nearest neighbours. It has recently been shown that long range coupling can be achieved through utilization of lateral leakage radiation in thin, shallow ridge silicon photonic waveguides [3]. This paper shows that the ability to break free from the limitations of simply connected waveguide topologies offers new opportunities for the realization of complex, multi-port quantum gates.

One of the most fundamental elements required for integrated optical devices, especially quantum devices, is the beamsplitter. This is the essential element for any interferometer, and can also be used (with trivial phase control) to effect a Hadamard rotation [4]. In the two-photon subspace, the beamsplitter can show the Hong–Ou–Mandel effect [5], one of the clearest non-trivial experiments to highlight the fundamental differences between classical and quantum optics. In integrated optics, a beamsplitter is typically realized through the use of a directional coupler. This is a device where two waveguides are brought into close proximity so that evanescent coupling causes population to tunnel between the waveguides. Truncating the device to the appropriate length then effects the desired beam splitting ratio. Although, in principle, it should be relatively easy to build such devices; in practice, any lack of control in the actual waveguide size leads to a lack of control in the evanescent tunnelling, and hence the length of device required to achieve a particular beam splitting ratio will be effectively unknown. A common solution is to post-select devices from a suite of similar devices, or alternatively, phase shifting elements such as heaters, can be used to fine tune and reconfigure devices [6].
Adiabatic passage promises a solution to issues of device variability that require postfabrication tuning. This is because adiabatic evolution goes as the ratio of tunnel matrix elements, rather than the absolute value of those elements. The tradeoff is that adiabatic devices are typically longer than their non-adiabatic counterparts, and whether the rewards of seeking an adiabatic versus non-adiabatic device are justified depends on the degree of device control required and the available footprint. Adiabatic methods for transport of population between states of the kind we envisage here began with the stimulated Raman adiabatic passage (STIRAP) protocol [7, 8], where robust transfer of population between atomic energy levels is effected by laser control, and the all-spatial variant that is sometimes called coherent tunnelling adiabatic passage (CTAP) [9–13]. There are also related schemes considering the adiabatic conversion of frequency modes of light [14].

Here we explore theoretically the potential for effecting geometric gates via long-range CTAP in thin, shallow ridge silicon-on-insulator (SOI) waveguides, shown schematically in figure 1(a). Specifically, we demonstrate several important gate designs including 50:50 beamsplitter and 1/3:2/3 beamsplitter, using a spatial version of the method outlined by Unanyan et al (USB) [15]. Our calculations are performed for both one and two photon input states. Further, we concatenate these devices to show an adiabatic non-deterministic controlled NOT (CNOT) gate, following the approach described by Ralph et al [38]. Our scheme utilizes a long-range common bus mode present in thin, shallow ridge waveguide devices [3]. This common bus mode provides a significant and new opportunity to develop planar geometries which are nonetheless not restricted to linear nearest-neighbour coupling. In this way, we see our approach as being more amenable to mass production, especially CMOS compatible fabrication, than truly three-dimensional approaches such as those described in, for example, [2, 16, 17].

Figure 1. (a) Schematic showing three thin shallow ridge SOI waveguides with a common bus (slab) mode. Coupling between waveguide and slab is effected by the overlap integral between the waveguide and the slab, and is therefore a function of the absolute position of the waveguide with respect to the slab. Each waveguide is assumed to be isolated from the other waveguides so that there is no appreciable evanescent tunnelling between the waveguides. (b) The tripod atom is the simplest system to realize USB style geometric adiabatic gates. Three ground states, |1⟩, |2⟩, and |3⟩, are coupled to a single excited state |0⟩ via optical fields with Rabi frequencies $\Omega_i$ for the transition between |i⟩ and |0⟩. The waveguide modes are equivalent to the ground states of the tripod atom, while the shared bus mode plays the role of the excited state.

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This paper is organized as follows. We first provide a brief introduction to thin, shallow ridge SOI waveguides, with emphasis on their effective refractive index and coupling to slab modes. As will be shown, control of both the magnitude and the sign of the coupling between waveguides and the slab can be achieved by the position of the waveguide relative to the slab mode. A change in the sign of the coupling leads to symmetry breaking mechanisms that are essential for USB-style geometric gates. With this understanding, we generate an effective Hamiltonian that can be used to effect arbitrary geometric gate sequences, and in particular we describe methodologies to realize an adiabatic power splitter, a Hadamard gate and 1/3:2/3 beamsplitter via the USB approach. After demonstrating one-qubit gates, we show the extension to two-photon gates, in particular showing that the well-known Hong–Ou–Mandel effect is preserved under conditions of adiabatic passage. Finally, we show the full state evolution for a non-deterministic linear optical controlled gate operating in the coincidence basis. This non-trivial two-photon, two-qubit entangling gate is completely simulated across eight optical modes (seven waveguide modes plus one bus mode), each of which can potentially have either 0, 1 or 2 photons.

2. Adiabatic evolution with thin, shallow ridge SOI waveguides

Thin, shallow ridge SOI waveguides when operating in TM polarization can exhibit lateral leakage behaviour [18]. Photons in a TM guided mode can leak into a lateral unguided TE slab mode, propagating at a specific angle to the waveguide axis due to polarization conversion at the ridge side-walls [19]. This TE slab mode can act as a bus mode to allow long-range communication between isolated waveguides [3, 20].

When the silicon slab is terminated, the continuum of TE slab radiation is discretised into discrete TE slab modes. On careful selection of the slab width, one of these TE slab modes can be phase-matched to the TM guided mode. Thus, the photons from the guided TM mode can couple to this TE slab mode. One particular method for varying the strength of the coupling between the guided TM mode and TE slab mode is by varying the relative location of the ridge waveguide on the slab [3, 20]. This technique opens up a new class of coupler that can enable interactions between multiple, well separated waveguides simultaneously, which is not possible in a traditional planar evanescent arrangement and has recently been proposed for CTAP devices [3]. While here we only consider interactions mediated via a single, discrete slab mode for clarity, generalizing our method so that coupling is via continuum states should be possible following the methods in [21, 22].

We consider a discrete silicon slab supporting a laterally defined bus mode |0⟩ with propagation constant $\beta_0$ and $N$ forward propagating waveguide modes |i⟩ of $\beta_i = k_0 n_i$, with effective index $n_i$, where $k_0 = 2\pi/\lambda_0$ is the free space wavevector for wavelength $\lambda_0$. Under these definitions, we may write down the system using a tight-binding Hamiltonian in
second quantized form as:

\[
H(z) = \beta_0 a_0^\dagger a_0 + \sum_{i=1}^{N} \beta_i a_i^\dagger a_i + \Omega_i a_i^\dagger a_i + \text{h.c.}, \tag{1}
\]

where \( a_i \) \((a_i^\dagger)\) is the photon annihilation (creation) operator acting on mode \( |i\rangle \) for \( i = 0 \ldots N \). Each waveguide is mutually isolated by separating an appropriate distance, ensuring there is no appreciable evanescent coupling, so that the waveguides only communicate through the common bus. The strength of this coupling \( \Omega \) is controlled by translating the waveguides laterally across the slab and this response is sinusoidal due to the nature of the bus mode \([3, 20]\). The coupling of a single waveguide to the bus is \( \Omega(z) = \Omega_{\text{max}} \sin[\beta x(z)] \), where \( x \) is the lateral waveguide location, assuming isolation occurs in the centre of the slab. The lateral waveguide dimension in turn varies as a function of the propagation dimension, \( z \), which is the mechanism to effect the adiabatic passage. The relationship between the lateral \( x \) and forward propagation \( z \) dimensions is controlled so that the couplings are varied adiabatically. The maximum coupling \( \Omega_{\text{max}} \) available depends on the waveguide dimensions \([19, 23]\) and can be calculated as the overlap integral between the bare bus TE and waveguide TM modes. This magnitude is represented as the imaginary effective index of a TM–TE coupled mode on an open slab \([19]\), or in the discrete case by observing the level of mode splitting throughout translation \([3]\). In the case where the slab width increases to accommodate additional waveguides, it is expected that this coupling will decrease as the maximal overlap of the single bus mode at any point decreases, resulting in longer devices. This scalability is of importance when considering more general Morris-Shore type devices \([24, 25]\).

In the discussion that follows we will adopt two separate notations. When we consider only the one-photon subspace, we will use the compact notation of defining the basis states by the position of the photon, i.e. we define \( |i\rangle \equiv a_i^\dagger |\emptyset\rangle \), where \( |\emptyset\rangle \) is the vacuum state of the system. However, when we deal with two-photon states, we will define the states by the occupation numbers of each mode, so for example the state \( |0110\rangle \equiv a_2^\dagger a_3^\dagger |\emptyset\rangle \). All of our simulations use a tight binding approach to solving the spatially varying Hamiltonian, and do not assume that the adiabatic limit is achieved.

Here we wish to inject light into port \([3]\) and arrive in an even superposition of both \(|1\rangle\) and \(|2\rangle\). The position of each waveguide is selected to provide ideal initial CTAP conditions \( |\Omega_i(0)\rangle = 0, |\Omega_1(0)\rangle = |\Omega_2(0)\rangle = |\Omega_{\text{max}}\rangle \). By translating the waveguides linearly across the slab, as illustrated in figure 2(a), the couplings are varied sinusoidally to effect the counter-intuitive pulse sequence. In particular, we have \( \Omega_3(z) = \Omega_{\text{max}} \sin \left[ \frac{\pi}{\Omegamax}(2z_{\text{max}}) \right] \) and \( \Omega_1(z) = \Omega_2(z) = \Omega_{\text{max}} \cos \left[ \frac{\pi}{\Omegamax}(2z_{\text{max}}) \right] \) shown in figure 2(b). The sin/cos coupling scheme has the nice property that the adiabaticity is constant throughout the protocol \([30–32]\). As \( \Omega_1 \) and \( \Omega_2 \) remain identical, an even power split arrives in each waveguide with the populations throughout the protocol described in figure 2(c). Combining the predicted paths, calculated population values and expected Gaussian mode profile of the waveguides gives a more visual representation of this transfer (figure 2(d)). This one input-two output device is equivalent to a Y-splitter but without the conventional restrictions of close proximity or ordering of the waveguides. There is therefore enhanced flexibility with the bus approach than more conventional approaches. This technique can also be extended to distribute population evenly across many waveguides in a method akin to that in \([17]\). The overall device length required to perform a successful adiabatic passage depends on \( \Omegamax \). The waveguide dimensions specified in \([3]\) provided a coupling length of 150 \( \mu \)m for a single waveguide and bus. It was shown that device lengths of \( \z_{\text{max}} \geq 2 \text{ mm} \) were required to successfully achieve robust adiabatic transport between two waveguides using the long-range bus, without significant population of the bus. This implies that to achieve successful CTAP behaviour the total device length must be at least longer than 15 coupling lengths. The magnitude of maximum coupling \( \Omegamax \) can be enhanced through waveguide engineering \([23, 33]\). Increasing the available coupling will decrease the coupling length and hence reduce the absolute device length \( \z_{\text{max}} \) required to maintain adiabaticity, although such optimization is not critical to explain our concepts.

3. Adiabatic power divider

Adiabatic techniques can be used for power division applications either by fractional adiabatic passage \([26, 27]\) or the use of additional waveguide modes \([28, 29]\). The presence of a shared bus offers an intriguing alternative technique. We first consider three waveguides acting as a tripod atom connected via a bus state (figure 1), the bus can be designed to be degenerate with the individual isolated waveguides, which will result in improved adiabaticity. All of our modelling was performed using the tight-binding Hamiltonian of equation (1) and was conducted in the adiabatic limit.

4. Controlled-ratio beamsplitters

The method of power splitting can be modified to effect robust quantum gates via the USB method \([15, 34]\). In this process, a double application of the power division is applied, with a change in the sign of the coupling applied between the first and second applications of the splitting. Because of the standing wave nature of the bus mode, the coupling between the waveguide and the slab varies sinusoidally with the waveguide position. The sinusoidal variation means that the sign of the coupling reverses in each standing-wave period. By ensuring that the forward and backward adiabatic passage crosses periods with the appropriate signs, the necessary symmetry breaking that is at the heart of the USB process can be achieved by waveguide translation alone.

In the one-photon subspace, the quantum gate is specified without loss of generality to act on qubit subspace of \(|1\rangle\) and
[2], with [3] as an auxiliary mode, explained below. We assume that the system is initialized in an arbitrary superposition $|\psi\rangle = \gamma_1 |1\rangle + \gamma_2 |2\rangle$, and for convenience we express the qubit in the dark/bright state basis

$$|D\rangle = \frac{\Omega_3 |2\rangle - \Omega_1 |1\rangle}{\sqrt{\Omega_1^2 + \Omega_3^2}}, \quad |B\rangle = \frac{\Omega_1 |1\rangle + \Omega_2 |2\rangle}{\sqrt{\Omega_1^2 + \Omega_2^2}}.$$  (2)

Note that provided the ratio of the couplings remains constant, the compositions of the dark and bright states will not change. Ideally $|D\rangle$ remains completely isolated from $|0\rangle$, while $|B\rangle$ can be influenced using CTAP. The counter-intuitive sequence transfers $|B\rangle \rightarrow |3\rangle$, instigating a phase reversal of $|B\rangle$ before returning it will alter the superposition, performing a rotation of the single qubit. The magnitude of this rotation is set by $\alpha = \Omega_3/\Omega_1$, and the net effect of a double application of the adiabatic passage is the gate:

$$G = \frac{1}{1 + \alpha^2} \begin{bmatrix} \alpha^2 - 1 & -2\alpha \\ -2\alpha & 1 - \alpha^2 \end{bmatrix}. \quad (3)$$

We illustrate the operation of this gate with three examples.

Preparing the system with all population initially in $|1\rangle$, the waveguides are translated to provide counter-intuitive transfer to $|3\rangle$, $\Omega_3$ then goes negative transferring back to $|1\rangle$ and $|2\rangle$ (figures 3(a) and (b)). Translating all waveguides linearly ($\alpha = 1$) results in complete population transfer to $|2\rangle$ (figure 3(c)). Performing this operation takes the system from $|\psi_0\rangle = \gamma_1 |1\rangle + \gamma_2 |2\rangle \rightarrow |\psi_\circ\rangle = \gamma_2 |1\rangle + \gamma_1 |2\rangle$, which is an X-gate: the quantum equivalent of a NOT gate. The trajectory taken by the population is also depicted on the sphere shown in figure 3(d). This qutrit representation displays only the real part of the state amplitudes [35–37].

Other gate operations occur when $\alpha \neq 1$. The available gates are confined to rotations in the canonical X–Z plane for the qubit defined across modes $|1\rangle$ and $|2\rangle$. If we define the trajectory of the state with greatest coupling to be linear, i.e. so that the coupling between this waveguide and the slab is sinusoidal and maximal, then it follows that the other must trace out a curved trajectory. So without loss of generality, assuming $\Omega_1 > \Omega_2$, we have

$$\Omega_1(z) = \Omega_{\text{max}} \sin \left( \beta_0 x_1(z) \right), \quad (4)$$

$$\Omega_2(z) = \Omega_{\text{max}} \sin \left( \beta_0 x_2(z) \right) = \alpha \Omega_1, \quad (5)$$

$$\therefore \ x_2(z) = \frac{\sin^{-1}(\alpha \Omega_1 / \Omega_{\text{max}})}{\beta_0}. \quad (6)$$

The Hadamard gate is a commonly used quantum information primitive, and is equivalent (up to phase) to a conventional beamsplitter. Preparing the system entirely in $|1\rangle$ a successful Hadamard operation will result in the state being transformed to $(1/\sqrt{2}) (|1\rangle + |2\rangle)$. Using (3), a value of $\alpha = \tan (\pi/8) \approx 0.4142$ provides this behaviour. The waveguide trajectories required to provide this value of $\alpha$ are shown in figure 3(a), with the resulting coupling scheme in figure 3(b). Evolving this system adiabatically results in the expected Hadamard operation as demonstrated in figures 3(e) and (f).

Another important beamsplitter is the 1/3:2/3 beamsplitter. Beamsplitters with this splitting ratio form the basis for non-deterministic linear optical quantum computing [38, 39]. A suitable two port 1/3:2/3 beam splitter is designed with $\alpha = 0.5176$ and operation is shown in figures 3(g) and (h).
5. Two-photon operation

There are few studies that have explicitly considered adiabatic passage of more than one particle [40–42], without invoking some mean-field or other effective treatment (as in for example [43, 44]). We are not aware of any previous works that consider adiabatic multi-particle gates such as we describe here, and hence some explanation of the two-photon gate operation is required.

The one-photon calculations described above are indistinguishable from the results that would be obtained for a purely classical modal analysis. Although the operation of the adiabatic gates on two-photon states is exactly what should be predicted from an equivalent conventional device, the microscopic details of how the adiabatic device achieves two-photon interactions are interesting and non-trivial.

For two photons across four modes (three waveguides + one bus mode), there are ten states that need to be considered. These, along with the couplings between the modes, are shown in figure 4. The states are: \(|0011\rangle\), \(|0101\rangle\), \(|0110\rangle\), \(|1001\rangle\), \(|1010\rangle\), \(|1100\rangle\), \(|0002\rangle\), \(|0020\rangle\), \(|0200\rangle\), and \(|2000\rangle\), where as before the most significant digit denotes the number of photons in the bus mode, and the subsequent digits refer to the number of photons in waveguides one to three.

Figure 3. Single photon arbitrary \(X-Z\) gate operations; (a) the paths of \(|1\rangle\) and \(|3\rangle\) are unchanged in each case, where the different gate operations are provided by altering the trajectory of \(|2\rangle\) which is represented as red (dashed), blue (dotted–dashed) and cyan (plus markers) for \(\alpha = 1, 0.4142\) and 0.5176 respectively, (b) the coupling of \(|3\rangle\) is now allowed to turn negative instigating a break in the symmetry of forward and backwards paths, (c) \(\alpha = 1\) operates as a NOT gate completely transferring population from \(|1\rangle\rightarrow|2\rangle\), (d) qutrit representation showing the forward (blue) and backwards (red) dotted paths taken through evolution, (e) and (f) \(\alpha = 0.4142\) conforms to a Hadamard operation creating a 50:50 superposition of \(|1\rangle\) and \(|2\rangle\), (g) and (h) \(\alpha = 0.5176\) creates a \(1/3:2/3\) beamsplitter.
vectors

\[
|D_1^{(2)}\rangle = \begin{pmatrix} \Omega_1^2 + \Omega_2^2 + \Omega_3^2 - \Omega_3 \\ - \Omega_1 \\ - \Omega_2 \\ - \Omega_3 \end{pmatrix} |0011\rangle
- \frac{\Omega_1}{\sqrt{2} \Omega_3} |0101\rangle - \frac{\Omega_1}{\sqrt{2} \Omega_2} |0110\rangle + |0200\rangle,
\]

\[
|D_2^{(2)}\rangle = \frac{1}{\Omega_2 \Omega_3} |0011\rangle
- \frac{\Omega_1}{\sqrt{2} \Omega_3} |0101\rangle - \frac{\Omega_1}{\sqrt{2} \Omega_2} |0110\rangle + |0200\rangle,
\]

\[
|D_3^{(2)}\rangle = -\frac{\Omega_3}{\sqrt{2} \Omega_3} |0011\rangle + \frac{\Omega_3^2}{\sqrt{2} \Omega_2 \Omega_3} |0101\rangle
- \frac{\Omega_3}{\sqrt{2} \Omega_1} |0110\rangle + |0020\rangle,
\]

\[
|D_4^{(2)}\rangle = \frac{1}{\Omega_3} |0011\rangle - \frac{\Omega_3}{\sqrt{2} \Omega_2} |0101\rangle
+ \frac{\Omega_3^2}{\sqrt{2} \Omega_1 \Omega_2} |0110\rangle + |0002\rangle.
\]

Figure 4. There are ten states involved in the two-photon, four mode Hamiltonian. This figure indicates the pertinent states and the strengths of the couplings between them.

With this four-dimensional null space, it is difficult to gain insight into the exact properties of the null state for any given problem, but there are some important features that can be gleaned. Firstly, notice that there is no overlap with states with a single photon in the bus mode. This is desirable as it minimizes sensitivity to loss from this mode. There is, however, potential overlap with the bus mode from \( |D_1^{(2)}\rangle \). We have numerically confirmed that providing the system is initialized in either \( |0011\rangle \), \( |0101\rangle \) or \( |0110\rangle \), there is no population in \( |2000\rangle \) (up to the limits of the adiabaticity of our numerical experiment, as seen for example in figure 5), indicating that \( |D_1^{(2)}\rangle \) is not populated during the adiabatic gate operation. The absence of population in the bus is important as it means that the adiabatic gate is indeed robust against loss from the bus mode.

Considering the 50:50 beamsplitter for the case of two indistinguishable input photons. The system is prepared in the state \( \ket{a_1 a_2} \ket{\emptyset} = |0110\rangle \). Using the same coupling scheme as in the one-photon case, namely

\[
\Omega_1 = \Omega \sin(\pi z / z_{\text{max}}),
\]

\[
\Omega_2 = \Omega \tan(\pi/8) \sin(\pi z / z_{\text{max}}),
\]

our results are shown in figure 5. Note the smooth, adiabatic evolution. In this case, the initial state \( |0110\rangle \) is transformed to an entangled state at the midpoint of the protocol, with non zero population in \( |0110\rangle, |0101\rangle, |0011\rangle, |0002\rangle, |0200\rangle \) and a very small contribution from \( |0200\rangle \). This evolution should be contrasted with the corresponding case from the one-photon Mandel response as measurement of the output ports will project the photons to be either both at waveguide 1, or both at waveguide 2.

6. CNOT operation

The CNOT gate is a fundamental entangling gate and popular choice as member of a universal gate set for quantum computing [45]. This is a two qubit gate, where the state of the target is flipped conditional on the state of the control qubit. One method to generate a scalable, but non-deterministic CNOT gate between individual photons is through combinations of linear optical elements (beamsplitters) [46]. Here we show the set of results when applying one particular implementation (the coincidence basis implementation) of a non-deterministic CNOT gate, based on 1/3:2/3 and 50:50 beamsplitters, following Ralph et al [38].

The canonical coincidence-based implementation requires six photonic modes, here encoded in the spatial modes available to the photons and shown in figure 6. Modes 1–3 correspond to the modes of the control. Mode 1 is the vacuum state for the control, mode 2 the control in the 1 state, and mode 3 the control in the 0 state. Mode 4 is the target 1 state, mode 5 the target 0 state and mode 6 the target vacuum state. There are then five full gate sequences,
which we denote $G_1$–$G_5$. $G_1$ and $G_5$ are 50:50 beam splitters, while $G_2$–$G_4$ are 1/3:2/3 beamsplitters. Our implementation requires eight modes, so in addition to the six modes discussed already, there is the bus mode (denoted by mode 0) that couples all of the modes via CTAP, and an auxiliary mode (mode 7) which plays the same role as the auxiliary mode for the one photon gate. Each gate works using the methods described above, with population adiabatically transferred from the interacting modes and auxiliary mode, via the bus mode. In the canonical CNOT gate, the 1/3:2/3 gates are performed in parallel, however in our case, due to the shared bus and auxiliary modes, all gates must be performed sequentially.

The results of performing the full CNOT gate operation on the appropriate starting states is shown in figure 7. The total state space for the problem with eight modes and up to two photons per mode has dimensionality 6561. However only 64 states actually participate in the problem, and in the adiabatic limit, only 49 of these will have non-zero population. Nevertheless, we do not label all of the occupied modes in the evolution shown in figure 7, instead only highlighting the starting states, with the final states given in table 1. The various output configurations are labelled as success or failure on the basis of whether they correspond to heralded success or failure of the non-deterministic gate. As expected, the table shows that the adiabatic passage CNOT gate operates in the same way as a conventional, coincidence-basis CNOT gate [38], with the correct state appearing with probability 1/9.

7. Experimental considerations

We now turn our attention to some practical considerations around the manufacture of adiabatic gates via the methods outlined here. While we cannot treat every potential error mode, two of the most important are: photon loss from scattering off roughness in the slab and misalignment of the waveguides with respect to the slab mode. Photon loss is suppressed by the adiabatic operation of the gate [28], so here we briefly describe the effect of waveguide misalignment during fabrication. Here we may observe that our scheme is less robust than laser-defined protocols such as STIRAP, because in our case the tunnel matrix elements go like the sine of the position. Note that this also means that the bus approach is more robust than conventional waveguide adiabatic passage, where the tunnel matrix elements go as the exponential of the waveguide separation.
If we consider a silicon on insulator device operating at a free-space wavelength of 1.5 μm, fabricated via e-beam lithography, then the relative error between waveguides is likely to be insignificant. However the alignment of the waveguides with respect to the slab is limited by the fabrication accuracy, which is likely to be of order nm. From this viewpoint, we consider a lithographically-defined waveguide protocol to effect a particular gate, but an error that misaligns the waveguides in one axis, by an amount \( \epsilon \). This error channel is potentially dangerous as the misalignment means that the ratio of the tunnel matrix elements, which is required for the adiabatic operation, will not be constant throughout the protocol. For a set of designed waveguide trajectories, \( x_1(z), \ x_2(z) \) and \( x_3(z) \), which are set to produce an ideal Hadamard operation (i.e. \( \alpha = \tan (\pi/8) \)), we have recalculated the gate fidelity, where the tunnel matrix elements have been replaced by

\[
\Omega_1'(z) = \Omega \sin \left\{ \frac{\pi [x_1(z) + \epsilon]}{z_{\text{max}}} \right\},
\]

\[
\Omega_2'(z) = \Omega \sin \left\{ \frac{\pi [x_2(z) + \epsilon]}{z_{\text{max}}} \right\},
\]

\[
\Omega_3'(z) = \Omega \sin \left\{ \frac{\pi [x_3(z) + \epsilon]}{z_{\text{max}}} \right\}.
\]

Under these conditions, with \( z_{\text{max}} = 1 \text{ mm} \), we retrieve the gate fidelity shown in figure 8. It is important to observe that the error is a smooth function of the misalignment, as is expected for an adiabatic process, and that fidelities of greater than 99.2% are achieved for \( |\epsilon| < 10 \text{ nm} \).

**Table 1.** Truth table/output modes for the adiabatic CNOT gate operation. State definitions are in the text, and designation specifies whether the output state is a failure mode or heralded success.

<table>
<thead>
<tr>
<th>Input state/configuration</th>
<th>Output configuration</th>
<th>Probability</th>
<th>Designation</th>
</tr>
</thead>
<tbody>
<tr>
<td>C_0T_0-{00010100}</td>
<td>00000110</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00001010</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00001100</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00002000</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00010010</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00010100</td>
<td>1/9</td>
<td>Success</td>
</tr>
<tr>
<td></td>
<td>00011000</td>
<td>1/9</td>
<td>Success</td>
</tr>
<tr>
<td></td>
<td>00020000</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td>C_0T_1-{00011000}</td>
<td>00000110</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00002000</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00001010</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00001100</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>01000010</td>
<td>2/9</td>
<td>Failure</td>
</tr>
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<td>01010000</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>01100000</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td>C_1T_0-{00100100}</td>
<td>00100010</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00110000</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>01000010</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>01010000</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td>C_1T_1-{00101000}</td>
<td>00100010</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>00110000</td>
<td>1/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>01000010</td>
<td>2/9</td>
<td>Failure</td>
</tr>
<tr>
<td></td>
<td>01010000</td>
<td>2/9</td>
<td>Failure</td>
</tr>
</tbody>
</table>

**8. Conclusions**

We have shown several designs based on adiabatic long-range couplers that can be useful in the distribution of power within integrated photonic circuits and have also demonstrated how this concept can be extended to quantum information processing specifically in the form of Hadamard and NOT gates, 1/3:2/3 splitters and describe how to perform an arbitrary X-Z rotation on a photonic qubit. We extend this to demonstrate a two-photon Hong–Ou–Mandel effect, and show the design of a complete non-deterministic linear optical CNOT gate. The feasibility to realize these designs in a planar CMOS compatible platform is very attractive in regards to large scale integration, fabrication accuracy and circuit complexity. While this work focusses on thin, shallow ridge SOI waveguides, it can be applied to other high index contrast materials that exhibit an accessible long-range bus mode. This technique may be useful in realising more general Morris-Shore type devices where large scale integration is of interest. As discussed, this scalability can require an increase in the overall device length to accommodate additional waveguides, as the overlap between each waveguide to a single bus mode decreases with increased slab width. The adiabatic nature of
these devices results in robust and repeatable signal transfer which is insensitive to variations in the device length.

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Chapter 5

Conclusions

This thesis has introduced a new technique to overcome the limits of nearest neighbour interactions which is suitable for mass-manufacture in a standard integrated silicon photonics foundry. The platform is based on adiabatic transfer and therefore is robust, this will make it a promising platform for sensitive quantum information experiments.

The thesis adapted lateral leakage ridge waveguides and showed that they could be fabricated using a standard silicon photonic foundry. The primary contribution of this thesis was a new concept, combining the leakage behaviour of these ridge waveguides with the adiabatic transfer concepts of CTAP and STIRAP to create a flexible long-range and arbitrarily configurable adiabatic photonic network. The potential of this platform was demonstrated by simulating how several important gates could be realised which, due to the novel photonic bus, could exceed the capabilities of previous planar photonic structures based on nearest neighbour interactions. The ultimate goal of a functioning quantum computer is still a future technology, however the knowledge base extended in this study will hopefully progress us a few steps closer.

5.1 Research summary

The primary outcome of this work was in the development of an integrated photonic platform that could break the restrictions of nearest neighbour interactions.

Chapter 2 was concerned with the design, fabrication and characterisation of shallow etched waveguides of varying widths and lengths, over a range of operating wavelengths. The direct imaging and quantifications of the lateral TE radiation was obtained for the first time in the literature. The high fabrication accuracy offered by silicon photonics foundries provided devices matching well to simulated predictions. This predictability increases the confidence in developing more
complicated structures that can make use of the lateral leakage effect.

Chapter 3 investigated the feasibility of designing an optical network that uses the lateral leakage mechanism in shallow ridge waveguides to overcome the limitation of nearest neighbour interactions. The lateral leakage acts as a coupling mechanism that allows pairs of waveguides to communicate, even with large waveguide separations, through the leakage acting as a ‘bus mode’ to spatially transfer photonic information. The application of CTAP provides deliberate transfer between waveguides while limiting the excitation of the laterally leaking mode, which decreases the propagation losses significantly. This CTAP protocol is robust to fabrication variations and in ideal situations will perform the transfer without any radiation. This high fidelity signal transfer is imperative in quantum information systems where it is necessary to transfer fragile qubits between functional elements. The coupling between each waveguide and bus mode is controlled by their relative lateral positioning on the silicon slab to provide effective CTAP population transfer. Full vectorial mode matching techniques showed that it is possible to transmit light using the bus mode as a mechanism for connecting the waveguides, while simultaneously limiting the amount of power that will be lost due to non-ideal scattering of the radiation. This shows an important new feature in waveguide transfer that is unprecedented and can completely bypass waveguides positioned between the target and source waveguides. The consequence of this adiabatic technique is that devices need to be longer than the equivalent directional couplers, with devices suggested on the order of 1 mm. The relative size of the devices can be further reduced by applying enhanced lateral leakage engineering such as providing a deeper etch, or higher order ridge structures.

It is important to note, that this optical bus design provides features that have not been possible in planar photonics in the past, breaking the restrictions placed on the relative waveguide geometries. As the waveguides are communicating via a common bus, this enables the interaction between any pair of waveguides in a potential $N \times N$ arrangement. This study provides the foundation for advanced low loss couplers with novel topologies suitable for the transport and manipulation of qubits. Such devices are conceived in the final chapter.

Chapter 4 includes a discussion on how to perform quantum information operations using these new lateral leakage CTAP elements. More specifically how to affect two-input, two-output beam splitting operations in a quantum setting. Such components are key in the design of new computing technologies based on light instead of electronics. Tight binding simulations show how to affect a variety of important operations by stacking functional CTAP elements in series. This included tracking a full 6561 dimensional Hilbert space when calculating the full two-photon CNOT gate, which hints at the power of just one single integrated quantum device. The use of the CTAP protocol also provides some protection against variations in the bus mode coupling.
5.2 Future outlook

I believe that this new photonic platform will present fertile ground for future researchers, particularly experimentalists pursuing integrated quantum information processing systems where this platform may provide unique opportunities.

The lateral leakage waveguides of Chapter 2 could benefit from additional characterisation, particularly in recording the effective index as a function of various waveguide parameters. Control over the effective index can control the particular angle of TE radiation. Analysing the purity of the radiated TE mode and how this depends on the radiation aperture will also be important for applications that require an integrated photonic antenna.

The most obvious next research goal will be in the fabrication and testing of the novel silicon photonic CTAP elements introduced in Chapter 3. However, there are some design considerations to make for the next fabrication. The design in Chapter 3 utilised the translation of waveguides across a discrete slab to provide the expected counter-intuitive coupling conditions. In an actual fabricated chip there will be a variation in the sidewalls that will decrease the robustness of the protocol. Work could be done to verify if the adiabatic passage technique is valid in an open boundary arrangement operating via a continuum, instead of a discrete bus mode. This will require a different method for controlling the coupling to the bus than is suggested by translating the waveguides in Chapter 3. Methods are available to achieve this including changing the waveguide widths, or providing an active tunable overlayer.

The work outlined within this thesis can be advanced by extending the gating technique introduced in Chapter 4 to more general scalable schemes such as Morris-Shore type devices. Implementing a Toffoli gate design would be of importance due to it being a key ingredient of quantum error correction systems. New applications can be explored such as the introduction of a silicon photonic bus to reduce the complexity in routing design within on-chip high-speed photonic links. Implementation of a dynamically reconfigurable switching network will require active switching with either electronic, or optical signals. As only one polarisation is affected by lateral leakage this will allow the design of polarisation splitters. The resonant lateral leakage may be further exploited by exposing the waveguide to an external measurand, which may find uses in bio-sensing applications.
Bibliography


