

Abstract

Advances in Silicon Photonics: Integrated Laser Stabilization, Isolation and Applications

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Silicon photonics has revolutionized integrated optical systems, offering compact, scalable, and high-performance solutions for applications in telecommunications, metrology, and quantum technologies. However, key challenges remain in achieving stable, high-performance on-chip lasers due to excess phase noise, frequency instability, and sensitivity to unwanted optical feedback. The absence of CMOS-compatible optical isolators and ultra-stable frequency references further limits the robustness of integrated photonic systems. This dissertation presents novel approaches to address these challenges by developing integrated laser stabilization techniques and broadband optical isolation solutions, leveraging micro-Fabry-Perot (μ FP) reference cavities and acousto-optic interactions.

A crucial requirement for high-performance integrated photonics is a stable and narrow-linewidth laser source. This work develops an approach for co-packaging ultra-high-quality factor, low-noise μ FP cavities with photonic integrated circuits to enable advanced laser stabilization techniques such as self-injection locking (SIL) and Pound-Drever-Hall locking directly on-chip. The μ FP cavities used in this study are fabricated via wafer-scale processing, resulting in vacuum-sealed resonators with sub-milliliter volumes. These cavities provide exceptional frequency stability with minimal thermal noise, making them ideal for integrated photonic applications.

To efficiently interface μ FP cavities with on-chip lasers, this work introduces a novel reflection transformation circuit (RTC) that enables stable optical coupling. A significant challenge in using Fabry-Perot cavities for laser stabilization is their high-

reflectivity response, which can destabilize laser operation. The RTC overcomes this issue by reshaping the reflection spectrum, allowing for stable SIL and PDH locking without optical circulators. Using this approach, we demonstrate a SIL laser with a record-low phase noise (-97 dBc/Hz at 10 kHz offset) and fractional frequency stability of 5×10^{-13} at 10 ms in an integrated silicon photonic system. These results establish a path toward ultra-stable, narrow-linewidth semiconductor lasers for applications such as high-precision spectroscopy, optical frequency metrology, and atomic clocks.

Beyond laser stabilization, this dissertation addresses the critical challenge of on-chip optical isolation. Traditional isolators rely on magneto-optic materials that are incompatible with CMOS processing, limiting their integration with photonic circuits. The lack of a practical on-chip isolator reduces the scalability of integrated photonic systems, as unwanted optical feedback can degrade laser coherence and introduce noise. To overcome this limitation, we develop a terahertz-bandwidth nonmagnetic optical isolator based on acousto-optic interactions in silicon photonic waveguides. By leveraging electrically driven phonon scattering, this approach enables broadband optical isolation without requiring magneto-optic materials. Through multimode waveguide dispersion engineering, we demonstrate an isolation bandwidth exceeding 2 THz with insertion losses as low as -2 dB. This level of performance surpasses existing integrated optical isolation techniques, offering a practical solution for protecting on-chip lasers from destabilizing back-reflections.

These innovations in laser stabilization and optical isolation represent a significant advancement in silicon photonics, addressing fundamental limitations in integrated photonic systems. The co-packaging of μ FP reference cavities with photonic circuits provides a new level of stability and coherence for integrated laser sources, enabling ultra-low-noise microwave photonics, coherent optical communications, and high-precision sensing. Meanwhile, the development of a CMOS-compatible, broadband optical isolator enhances the robustness of integrated lasers and amplifiers,

making fully integrated photonic systems more scalable and reliable.

Beyond telecommunications and metrology, the technologies demonstrated in this work have broader implications for emerging fields such as quantum photonics, where coherence and stability are essential for quantum computing and sensing. The ability to integrate ultra-stable laser sources and nonreciprocal optical components on a single chip paves the way for miniaturized, portable precision photonic systems with applications in navigation, remote sensing, and imaging.

Advances in Silicon Photonics: Integrated Laser
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Contents

Acknowledgments	xv
1 Introduction	1
1.1 Silicon Photonics	1
1.2 Integrated Laser and Amplifier Technologies	2
1.3 Integrated Laser Stabilization Techniques	4
1.3.1 Pound-Drever-Hall(PDH) Locking	5
1.3.2 Self-Injection Locking	6
1.3.3 Reference Cavities	7
1.4 Integrated Laser Protection	8
1.4.1 Laser Destabilization from Unwanted Feedback	8
1.4.2 On-chip Nonmagnetic Isolator	9
1.5 Summary of Dissertation	10
2 Integrated Silicon Nitride Reference Cavity	13
2.1 Ultra-low-loss Silicon Nitride Waveguide	13
2.2 Roughness-Induced Scattering Loss	15
2.2.1 Volume current method	15
2.3 Scattering Loss Estimation for Dilute Si ₃ N ₄ Waveguide	21
2.4 Conclusion	23

3	Co-Packaged micro-Fabry Perot Cavities for On-Chip Laser Stabi-	24
	lization	
3.1	Introduction	24
3.2	Co-integration of μ FP Cavities	27
3.3	Reflection Transformation Circuit for PDH Locking	29
3.3.1	System concept of on-chip poor man's isolator	29
3.3.2	Inverse-design optimized polarization splitting grating coupler	33
3.3.3	Experimental results	35
3.3.4	Co-packaged μ FP cavity with poor man's isolator interface through PSGC	38
3.4	Reflection Transformation Circuit for Self-injection Locking	40
3.4.1	System concept	40
3.4.2	SIL laser with co-packaged μ FP cavity	43
3.5	Conclusion	49
4	On-chip Terahertz Bandwidth Nonmagnetic Isolators	52
4.1	Introduction	52
4.2	System Concept	54
4.2.1	Electrically driven acousto-optic scattering	56
4.2.2	Cavity enhanced acoustic-optical scattering	59
4.2.3	Extending isolation bandwidth	60
4.3	Experimental Results	68
4.3.1	Terahertz bandwidth isolation based on dispersion balanced multimode waveguide	68
4.3.2	Wavelength tunable wideband isolation based on length bal- anced singlemode waveguide pair	74
4.4	Discussion and conclusion	77

5	Conclusion	80
5.1	Summary	80
5.2	On-going projects and outlook	83

List of Figures

2.1	Top view of the dielectric distribution for (a) a realistic waveguide and (b) an ideal waveguide.	15
2.2	(a) Illustration of a perturbation at the dielectric boundary. (b) Comparison between the ideal and perturbed boundaries, highlighting the importance of the permittivity choice at the blue point. (c) Green's function calculation for a dipole near a semi-infinite dielectric boundary with the dipole oriented parallel to the interface. (d) Green's function calculation for a dipole near a semi-infinite dielectric boundary with the dipole oriented perpendicular to the interface.	17
2.3	(a, i) Schematic of a dilute Silicon Nitride waveguide. (a, ii) Fundamental TE optical mode of the waveguide. (a, iii) Typical surface roughness profile of the Silicon Nitride core. (b) Comparison of the scattering loss estimation to experimental results. (c) Scattering loss prediction for various core geometries. (d) Scattering loss comparison of TE and TM mode.	22

3.1	a , Survey of resonator thermal noise limit at 10 kHz and quality factor for on-chip ring cavity, WGM cavity and vacuum FP cavity. Ring resonator data is from [12, 32, 42–44], WGM resonator data is from [34, 45, 46], and vacuum FP resonator data is from this work and [35, 47]. Note that our cavity has $> 10^4$ times smaller form factor than [35] and doesn't require vacuum enclosure. b , i , Schematic illustration of the in-vacuum bonded μ FP cavity. ROC: radius of curvature. ii Photo of μ FP cavity. c , Optical ring-down measurement of μ FP cavity yielding quality factor of 4.26 billion. d Schematic of the integration architecture. A reflection transformation circuit is required to transform mirror-like reflection to the response desired by application. e Schematic of mode-matching approach from PIC to μ FP with i end-fire and ii grating coupler.	27
3.2	a , Depiction of the traditional setup of a poor man's isolator mechanism in free space, where PBS denotes a Polarizing Beam Splitter. b , Illustration of the circuit interface, which is modeled after the poor man's isolator, that incorporates a the Polarization Splitting Grating Coupler (PSGC). PIC: photonic integrated circuits. c , Schematic of PSGC. Light input from Arm 1 will be scattered into p-polarized light in free space and that from Arm 2 will be scattered into s-polarized light in free space. d Transfer matrix schematic of the system. . . .	30
3.3	Edge-coupler based on-chip poor man's isolator interface circuit. . .	31

3.4	a , Side view: The layered structure of the PSGC is presented, with a single mode optical fiber positioned perpendicularly to the PSGC. Top view: SEM image of the PSGC. Pattern is symmetrical over the white line. b , FDTD simulation showcases the scattered field 3um above the Si device layer, prompted by an incident waveguide mode from bottom port. c , Empirical measurement of the coupling efficiency and crosstalk of the PSGC. Efficiency is characterized as the transmission from Arm 1 to the fiber, while crosstalk is defined as the transfer from Arm 1 to Arm 2. d , Zoom-in of the crosstalk data presented in (c). e , Scatter plot showing the performance among ten distinct PSGCs fabricated on the same chip.	34
3.5	a , Measurement setup for the characterization of the reflection cancellation circuit. DUT: device under test, PD: photo detector. b-c , Measurement results with fiber mirror (FM) and fiber Bragg grating cavity (FBGC) as DUT. d , μ FP array. e , Fiber GRIN lens collimator assembled to match μ FP cavity mode. f , Measurement results with fiber-GRIN- μ FP as DUT.	35
3.6	a , Schematic of the approach for co-integrating μ FP cavity with signal redirection and back-reflection isolation. b, i Cross section of co-integrated μ FP cavity. ii Optical picture for co-integrated module. c , Experimental data validating signal redirection and characterizing back-reflection isolation ratio. d , Outlook for on-chip PDH locking system. PM: phase modulator. PD: photo detector. LD: laser diode. EIC: electric integrated circuits.	39

- 3.7 **a**, Schematic illustration of on-chip circuit design enabling self-injection locking to co-integrated μ FP cavity. **b**, Photograph of DFB laser self-injection locked to μ FP cavity through interface chip. **c**, Schematic of experimental setup for noise measurement. Black lines indicates electrical control, and red lines indicates optical circuits. PD: photo detector. ESA: electrical spectrum analyzer. AOM: acousto-optical modulator. The output of SIL laser either beats with another stable laser or performs self-heterodyne with 800 m delay line for the ESA to measure phase noise. **d**, Phase noise spectrum for SIL and free running laser. Free running laser phase noise is measured using self-heterodyne. The faded curve indicates the noise is limited by the measurement noise floor. **e**, RF beat note spectrum of the SIL laser and a stable laser reveals linewidth compression when locked. **ii** is the zoom in for **i**. Free running beat note is taken with 91 kHz RBW, and locked beat note is taken with 3 kHz RBW. **f**, Allan deviation of the SIL laser measured from frequency counter with and without linear drift removal. 41
- 3.8 **a**, Schematic of the system. **b**, Response curve for a DFB laser directly coupled to a μ FP cavity. **c**, Response curve for a DFB laser directly SIL to a μ FP cavity with perfect mirror reflection removal, shown for various feedback phases. **d**, Noise reduction factor for a DFB laser directly SIL to a μ FP cavity with perfect mirror reflection removal as a function of the feedback phase. **e**, Response curve for a DFB laser directly SIL to a μ FP cavity with imperfect mirror reflection removal. Cases (i–iii) correspond to $\phi_{m0} - 2\phi_{c0} = 0, -\pi/20, -\pi/10$ 45

3.9	a , RF beat note spectrum of the SIL laser and a stable laser with a sweeping current on the DFB laser yields a 265 MHz locking range. b , Theoretical locking curve. The SIL laser entering and exiting the locked state are due to relaxation oscillation of the DFB laser. Off-resonance curve deviates from free running due to the residual broadband mirror reflection.	46
3.10	Collection of the previous results shown in [37, 38, 47, 58, 64, 65]. . . .	50

4.1 **a** Dual-rail representation of the bandwidth-scalable isolator for forward (left) and backward (right) operation. AOM1 and AOM2 denote acousto-optic modulators 1 and 2, respectively, acting as nonreciprocal beam splitters. **b** The nonreciprocal acousto-optic modulator in our system partially scatters the forward-propagating symmetric mode at frequency ω_1 into an antisymmetric mode at frequency $\omega_2 = \omega_1 + \Omega_0$, while partially scattering the backward-traveling antisymmetric mode at frequency ω_2 into a symmetric mode at frequency $\omega_3 = \omega_2 + \Omega_0$. Here Ω_0 represents phonon angular frequency. A mode-multiplexer (MM) is used to selectively excite and extract different spatial mode in this multimode waveguide. For detailed discussion on AOM nonreciprocity, please refer to SI. Sec I. **c** Theoretical response of the nonreciprocal multimode interferometer. Similar to a conventional Mach-Zehnder interferometer, the response exhibits a cosine-shaped dependence on optical frequency. The nonreciprocal beam splitter—implemented here by the AOM—shifts the forward and backward responses, thus functioning as an isolator. By balancing the two arms of the interferometer, characterized by $n_{g,+}L$ and $n_{g,-}L$, the isolation bandwidth is increased. **d** Two complementary approaches for bandwidth-scalable optical isolators. Optical paths in the delay region for symmetric mode and anti-symmetric mode are denoted by solid and dashed line respectively, same as panel **a**. **i** Balancing the interferometer by transiting the regular multi-mode waveguide to a dispersion balanced ($n_{g,+} \cong n_{g,-}$) multimode waveguide. **ii** Balancing the intereferometer by mapping two modes into two identical waveguides ($L_+ \cong L_-$) through mode multiplexers (M3, M4). 55

4.2	a , Schematic of single section electro-optomechanical scattering. b , Dispersion diagram of optomechanical scattering. c , Single section measurement of three different devices with a RF drive at 2.36 GHz. i , ii refer to dispersion-balanced multimode waveguide device A and B iii refers to length balanced singlemode waveguide device.	57
4.3	a , Acoustic cavity simulation for the mode favorable for inter-mode scattering. b , Band diagram of phononic crystal. c , Dimensions for the multimode waveguide of AOM active region. d , S11 measurement of the IDT.	59
4.4	Illustration of the contribution from different group delay terms using Device B fitted data.	67
4.5	a , The isolator microscopy and experimental setup for characterizing the forward and backward operation of our isolator. SG: signal generator; Amp: RF amplifier; Splitter: 50/50 power splitter; PS: RF phase shifter; AOM: acousto-optic modulator; PD: photodiode; RFSA: RF spectral analyzer; M1/2: mode-multiplexer. Pink lines represent optical connections, blue lines represent high frequency (~ 2 GHz) electrical connections and green lines represent low frequency (55MHz) electrical connections. b , The on-chip acousto-optic modulator (AOM) comprises a piezoelectric interdigital transducer (IDT) fabricated on AlN and an optomechanical waveguide crafted from a suspended single silicon ridge. c , The dispersion-balanced waveguide is an unsuspended two-ridge waveguide. A SEM image is shown in ii and the designed dimensions of the two-ridge waveguide are labeled. Such waveguide provides similar group index of the symmetric (iii) and antisymmetric (iv) modes.	69

4.6 **a**, We measure the mode conversion efficiency as a function of the microwave drive frequency while driving only AOM 1 (dark blue in **i**) and driving only AOM 2 (light blue in **i**), which demonstrates almost identical responses. When both AOMs are simultaneously driven (**i**), the forward (backward) operation of our device equates to the sum (difference) of the responses from the two AOMs, leading to the high isolation contrast exhibited by our device. **b**, The insertion loss decreases as we increase the microwave drive power, approaching -2 dB. The maximum RF power is limited by the power handling of the device. **c**, The transmission of Device A operating in forward (red) and backward (green) directions, which demonstrates the 10 dB isolation bandwidth of 5.8 nm (725 GHz, shaded blue window in **i**, enlarged in **ii**). Over the isolation bandwidth, the device has -2.16 dB insertion loss and 12.4 dB peak isolation ratio. Measurements deviate slightly from theoretical predictions (black) due to pump depletion and intermodal crosstalk in the passive photonic components of the device. Note that higher contrast is seen between forward and backward transmission in (**a**) because these frequency-resolved measurements display only the frequency-shifted components of the transmitted light field. **d**, The transmission of Device B operating in forward (red) and backward (green) directions, which demonstrates the 10 dB isolation bandwidth of 16 nm (2 THz, shaded blue window in **i**, enlarged in **ii**). Over the isolation bandwidth, the device has -8.09 dB insertion loss and 28.1 dB peak isolation ratio. 71

4.7	a , The second dispersion engineering method incorporates two extra mode multiplexers to split the two waveguide modes into two single-mode waveguides. b , An optical micrograph of the fabricated device. The two-waveguide delay line is configured in a spiral shape, taking advantage of its robust dispersion properties that can withstand waveguide bends without significant degradation. c , Conversion efficiency spectrum with difference phase difference between two AOMs. Our device exhibits wavelength tunability (i-iii) and reconfigurability of the isolation direction (iii and iv), which can be adjusted by controlling the relative phase difference θ between the microwave drives for AOM 1 and AOM 2.	74
4.8	Review of on-chip non-magnetic optical isolators. Collection of the previous results shown in [22–27, 82–84].	77
5.1	a On-chip PDH locking circuit design. PM: phase modulator, TC: tunable coupler, PD: photo detector, PMI: poor man’s isolator we discussed in Sec 3.3.1. b Dual PDH locking for mm-wave generation. c Detailed GDS design of on-chip PDH locking chip. d Detailed GDS design of on-chip dual-PDH locking chip. e Zoom-in of the sub-mount and package for dual-PDH chip. f Zoom-out of the package for dual-PDH chip.	81
5.2	a On-chip dual-SIL circuit design. PS: phase shifter, TC: tunable coupler, PD: photo detector. b Zoom-in of the package design, with two DFB lasers sitting on two separate TECs. c Detailed GDS design of on-chip dual-SIL chip.	83
5.3	a Schematic of hybrid electro-SIL laser. BPD: balanced photo detector b Detailed GDS design of hybrid electro-SIL laser chip.	85

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Chapter 1

Introduction

1.1 Silicon Photonics

Integrated photonics has revolutionized the field of optical communications, signal processing, and quantum optics by enabling the integration of optical components onto chip-scale platforms. Among the various approaches, Silicon photonics stands out due to its compatibility with established complementary metal-oxide-semiconductor (CMOS) manufacturing processes. By leveraging mature Silicon-based fabrication techniques, Silicon photonics offers the potential for the mass production of highly efficient, compact, and cost-effective optical devices. Silicon photonics not only enhances the scalability of optical systems but also facilitates innovative applications in areas such as high-speed communications, quantum computation, and sensing technologies [1–3].

From an optical perspective, Silicon offers several advantageous properties. Its high refractive index contrast with cladding materials, typically Silicon dioxide, enables strong optical confinement, allowing for the miniaturization of photonic devices. Additionally, Silicon's transparency in the near-infrared spectrum, particularly at 1.3 μm and 1.55 μm , makes it an ideal platform for optical communication and sensing

applications. Furthermore, the ability to dope Silicon provides added functionality. For example, doping enables the creation of modulators via the plasma dispersion effect [4], integrating both electronic and photonic functionalities into photonic integrated circuits (PICs). More recently, people also introduced low loss Si_3N_4 in Silicon photonic platform [5], bringing more chip-scale functionality.

However, Silicon photonics also faces certain limitations. One of the key challenges is the integration of efficient light sources. Due to Silicon’s indirect bandgap, it is not suitable as a direct light emitter, and requires the use of external light sources or hybrid integration of III-V semiconductor lasers. Additionally, achieving ultra-stable on-chip light sources is difficult due to thermal-related noise from dielectric waveguides, which can degrade the coherence and stability of the laser output. Moreover, optical isolators, which are essential for protecting laser sources from back reflections, typically rely on magneto-optical materials, such as Ce:YIG, which are not compatible with CMOS fabrication processes. Addressing these challenges is crucial to fully harnessing the potential of Silicon photonics. This dissertation explores techniques, including novel circuit designs and co-packaging approaches, that aim to overcome these limitations and drive the next generation of Silicon photonic systems.

1.2 Integrated Laser and Amplifier Technologies

The integration of laser sources and optical amplifiers into Silicon photonic platforms represents a pivotal advancement toward fully integrated photonic systems. While Silicon photonics has successfully leveraged the advantages of Silicon for passive electrical-optical components, incorporating active elements such as lasers and amplifiers poses significant challenges due to Silicon’s indirect bandgap and lattice mismatch to direct bandgap III-V materials. Recent developments in integrating III-V semiconductor materials with Silicon have opened new pathways for efficient

on-chip light generation and amplification.

One prominent approach is heterogeneous integration, wherein III-V materials are directly bonded onto Silicon substrates. The direct bonding technique allows for the combination of the superior optoelectronic properties of III-V materials—such as efficient light emission and amplification—with the mature manufacturing infrastructure of Silicon. This synergy has led to the development of high-performance, scalable, and cost-effective integrated photonic devices [6].

Alternatively, advanced epitaxial growth techniques have been employed to directly grow III-V materials on Silicon substrates. The primary challenge in this approach is the lattice mismatch and thermal expansion coefficient differences between III-V materials and Silicon, which can introduce defects and degrade material quality. By introducing lattice-matching buffer layers, these disparities can be mitigated, enabling the growth of high-quality III-V crystals directly on Silicon [7]. This monolithic integration approach holds the promise of simplifying fabrication processes and enhancing integration density, as it circumvents the need for bonding separate materials.

In addition to III-V material integration, the doping of Silicon-based materials with rare-earth elements has emerged as a viable strategy for on-chip optical amplification. Specifically, Erbium-doped Si_3N_4 waveguides have demonstrated the ability to provide optical gain within the telecommunications wavelength bands. The low-loss characteristics of Si_3N_4 , combined with the amplification properties of Erbium ions, enable the realization of integrated optical amplifiers compatible with Silicon photonic platforms [8]. This approach leverages existing fabrication techniques and materials, offering a pathway toward fully integrated amplification without the need for heterogeneous material integration.

Despite these advancements, several challenges impede the performance of integrated lasers and amplifiers when compared to their bulk or table-top counterparts.

A significant issue is the noise performance of integrated lasers. On-chip lasers often exhibit higher phase noise and broader linewidths due to factors such as increased susceptibility to thermal fluctuations, limited cavity lengths, and material imperfections inherent in the integration processes. Efficiently stabilizing integrated lasers remains an active area of research. Another critical challenge is the sensitivity to reflective feedback. Lasers and amplifiers are inherently sensitive to reflections due to the dynamics of the gain medium, where even minor back-reflections can induce intensity noise, mode hopping, or even destabilize the lasing action entirely. In conventional optical systems, optical isolators based on magneto-optic effects are employed to mitigate this issue. However, integrating such isolators into Silicon photonic platforms is problematic because magneto-optical materials are incompatible with CMOS fabrication processes. Developing CMOS-compatible isolation techniques is essential for the reliable operation of integrated lasers and amplifiers.

Addressing these challenges is crucial for unlocking the full potential of Silicon photonics in applications ranging from high-speed communications to quantum information processing. This dissertation delves into novel integration strategies, circuit designs, and material systems aimed at overcoming the limitations of current integrated laser and amplifier technologies. By advancing the understanding of these components and their interactions within integrated platforms, we aim to contribute to the development of next-generation photonic systems that offer enhanced functionality, improved performance, and greater scalability.

1.3 Integrated Laser Stabilization Techniques

In this section, we will explore techniques for stabilizing integrated lasers and introduce several types of promising reference cavities adaptable to integrated photonic platforms.

1.3.1 Pound-Drever-Hall(PDH) Locking

Pound-Drever-Hall (PDH) locking is a powerful and widely-used technique for stabilizing laser frequencies by locking them to the resonance of a high-quality-factor optical cavity. This method is critical in numerous applications requiring high-precision laser stabilization, such as spectroscopy, atomic clocks, gravitational wave detection, and high-resolution metrology. Developed by R. V. Pound, Ronald Drever, John Hall, and their collaborators in the 1980s, PDH locking allows the laser to achieve an extremely narrow linewidth and reduced frequency noise, making it highly suitable for modern photonic systems where high-level of precision and stabilization is required [9, 10].

The principle behind PDH locking involves modulating the phase of the laser beam and then detecting the reflected signal from the optical cavity. A phase modulator, such as an electro-optic modulator (EOM), is used to impose sidebands on the laser light at a specific modulation frequency. The modulated laser beam is then directed towards the high-quality-factor optical cavity. When the laser frequency is slightly off resonance, interference occurs between the carrier and sidebands, resulting in an error signal that is detected by a fast photodetector. This error signal carries information about the frequency difference between the laser and the cavity resonance. By analyzing this error signal, precise adjustments can be made to maintain the laser at the desired frequency using an electronic feedback loop that controls the laser's frequency actuator, such as a piezoelectric transducer (PZT) or the laser current.

This feedback-based technique significantly mitigates noise and drift in the laser frequency, ensuring stable operation even in challenging environments. One of the key advantages of PDH locking is its ability to achieve a steep and highly sensitive error signal, allowing for precise corrections to be applied to the laser frequency. This ensures that the laser remains locked to the cavity resonance, even in the presence of external perturbations such as mechanical vibrations or temperature fluctuations. Additionally, PDH locking can be implemented with relatively low-cost components,

making it accessible for a wide range of laboratory and industrial applications.

PDH locking has become an essential tool in many advanced research fields. For instance, in gravitational wave detectors like LIGO, PDH locking is employed to stabilize the laser systems that serve as the core of the interferometric measurement. In high-resolution spectroscopy, the technique enables the precise control of laser frequency, allowing researchers to probe atomic and molecular transitions with exceptional accuracy. Similarly, in optical atomic clocks, PDH locking is used to stabilize the clock lasers to an ultra-stable reference cavity, contributing to the unprecedented accuracy of modern timekeeping systems. The versatility and robustness of PDH locking make it an indispensable method for any application requiring stable, narrow-linewidth laser sources.

1.3.2 Self-Injection Locking

While PDH locking is well-known for its ability to achieve robust stabilization by locking a laser to the resonance of an optical cavity, it typically requires relatively complex electronic modulation and feedback systems. In contrast, self-injection locking (SIL) provides a simpler approach by utilizing external resonant optical feedback to stabilize the laser. This makes SIL particularly attractive for compact, chip-scale systems where reducing complexity and component count is essential. Both methods can achieve ultra-narrow linewidths and low frequency noise, but SIL's simplicity makes it highly suitable for integrated photonic applications, whereas PDH locking remains a preferred choice in scenarios demanding the highest possible precision and stability, such as atomic clocks and gravitational wave detectors.

The basic principle of self-injection locking involves reflecting a portion of the emitted laser light back into the gain medium after passing through a resonator. The resonator provides a selective frequency response, effectively locking the laser's frequency to one of its resonant modes. This feedback loop allows the laser to op-

erate with significantly enhanced stability compared to free-running conditions. The technique is particularly effective with integrated photonic platforms, where compact microresonators can be co-located with the laser on a single chip [11], enabling high performance while minimizing size, weight, and power consumption.

Self-injection locking has demonstrated considerable promise in integrated laser technologies, with recent advancements showing successful integration of high-Q reference cavities on silicon photonic chips [11,12]. This has opened pathways for practical implementation of ultra-low-noise lasers, which are key to advancing technologies in optical communication, lidar, and frequency comb generation. As the demand for high-performance lasers continues to grow, self-injection locking remains a focal point of research in the quest for stable and high-quality laser sources. Moreover, how to maintain a stable long term robust SIL as what PDH locking does is still remain to be studied.

1.3.3 Reference Cavities

As we discussed above, no matter which approach, a stable, high-quality factor optical reference cavity is needed. In this section, we will introduce several candidates of reference cavities with the potential of integration to Silicon photonics. Fabry-Perot(FP) cavity, formed by two parallel high reflectivity mirrors, is widely used as reference cavities. It's advantages to serve as reference cavity because no dielectric is required between two high reflectivity mirrors. It's very beneficial to leave the gap in vacuum. One is it can eliminate any loss inside of the dielectric, providing highest quality factor. The other is it can also eliminate thermal noise associated with dielectric materials. With optical field lives virtually in vacuum, such vacuum FP cavity can provide extraordinary stability. Traditionally, such vacuum cavities have bulky form factor and it also requires complicated vacuum enclosure to keep the vacuum. Recent advances provides wafer level fabrication approaches for miniaturized micro-

FP cavities [13], with wafer bonding techniques, it can seal vacuum inside, eliminate complicated vacuum enclosures. Such vacuum FP cavity with millimeter lever form factor makes it suitable to be integrated with PICs. However, due to the fact that the reflection spectrum of a FP cavity is typically a wide-band mirror like reflection. In partical use, it often comes with a circulator to remove the backreflection and redirect the signal. This issue prohibits the introduction of FP cavity into PIC platform for a long time, due to the factor that integrated circulator is not currently available. In this dissertation, we introduce novel approach to solve this problem and effectively utilize micro-FP cavity in integrated platform.

Another promising candidate is ultra-low-loss Silicon Nitride ring resonators [5]. The waveguide based structure makes it suitable for chip scale applications. Due to the fact that thermal noise associated with dielectrics is inversely proportional to optical mode volume, people use meters long spiral resonators to reduce the thermal noise limit on chip. Although the performance is still not comparable with vacuum FP cavities, it's suitable for integrated system with less tight requirement on noise. In this dissertation, we discuss the theoretical model for roughness induced scattering loss in such waveguide based resonators.

1.4 Integrated Laser Protection

1.4.1 Laser Destabilization from Unwanted Feedback

Lasers and amplifiers are highly susceptible to unwanted reflections and backscatter due to dynamics of the gain media. For example, when a laser receives uncontrolled feedback from external reflections, this feedback interferes with the gain dynamics of the laser cavity, leading to various destabilizing effects, such as increased intensity noise, mode hopping, frequency shifts, or in the worst-case scenario, complete loss of stable lasing [14]. This will be catastrophic for building robust ultra-stable on-chip

optical systems.

Unwanted feedback can originate from multiple sources, including interface reflections due to refractive index mismatches at waveguide terminations, roughness-induced back-scattering [15], or even reflections from optical components like photo-detectors and grating couplers. These reflections can couple back into the laser cavity and interact destructively or constructively with the circulating light, significantly impacting the phase and amplitude of the laser field. Such effects are particularly pronounced in integrated photonic circuits, where space and optical components are limited due to CMOS process limitation.

1.4.2 On-chip Nonmagnetic Isolator

The development of integrated optical isolators is a critical challenge in the field of silicon photonics. Optical isolators prevent unwanted light reflections from propagating back into a laser source, which can destabilize the laser and degrade system performance. Traditional optical isolators rely on magneto-optic effects, requiring the use of magnetic materials such as cerium-doped yttrium iron garnet (Ce:YIG). However, these materials are incompatible with complementary metal-oxide-semiconductor (CMOS) processes, making them impractical for integration into photonic circuits on silicon platforms.

Recently, time-modulated systems have been employed to create nonmagnetic forms of nonreciprocity [16–18], leading to impressive demonstrations of optical isolation using electro-optics [19, 20], acousto-optics [21], and Kerr nonlinearities [22]. However, current implementations are limited by constraints associated with the response of integrated resonators [22–25], time-modulation frequency [20], or the phase-matching bandwidth of traveling wave interactions [26, 27], resulting in relatively narrow (on the order of gigahertz) isolation bandwidths. While such narrowband isolators may be useful in specific applications, achieving wideband (greater than terahertz)

optical isolation remains critical for the integration of wideband optical amplifiers, swept lasers, and broadband light sources. Furthermore, wideband optical isolators eliminate distortion and instabilities caused by unwanted feedback, enabling scalable and frequency-agnostic integration of cascaded photonic components, which is key to unlocking the full potential of integrated photonics. In this dissertation, we introduce a novel approach to create low loss optical isolators with scalable bandwidth.

1.5 Summary of Dissertation

This dissertation explores key advancements in silicon photonics, focusing on integrated laser stabilization, broadband optical isolation, and novel photonic circuit interfaces. The research addresses fundamental challenges in achieving stable, narrow-linewidth lasers, mitigating unwanted optical feedback, and implementing CMOS-compatible optical isolators. The following chapters detail the innovations developed in this work.

In Chapter 2, we discuss the development of ultra-low-loss silicon nitride waveguides and their application in integrated photonic reference cavities. The fundamental limitations of silicon photonics, particularly thermal noise and waveguide scattering loss, are analyzed. A theoretical model based on the volume current method is introduced to quantify roughness-induced scattering losses in high-aspect-ratio silicon nitride waveguides. This model enables precise estimation of optical losses, providing guidelines for optimizing waveguide geometries to achieve high-quality-factor (Q) resonators. The chapter also discusses the trade-offs between different waveguide designs and polarization modes, demonstrating that ultra-low-loss silicon nitride cavities are a viable solution for laser stabilization applications.

In Chapter 3, we introduce the integration of vacuum-gap micro-Fabry-Perot (μ FP) reference cavities with photonic integrated circuits to achieve ultra-stable laser

operation. This chapter details a wafer-scale fabrication process for μ FP cavities, which enables high-Q resonators with vacuum-sealed enclosures to eliminate thermal noise. A critical challenge in integrating μ FP cavities with PICs is their strong mirror-like reflection, which can destabilize on-chip lasers. To address this, a novel reflection transformation circuit is introduced to reshape the cavity reflection spectrum, enabling stable self-injection locking and Pound-Drever-Hall locking without requiring optical circulators. Using this co-packaged μ FP cavity system, a record-low laser phase noise of -97 dBc/Hz at 10 kHz offset and a fractional frequency stability of 5×10^{-13} at 10 ms are demonstrated. These results represent a major advancement in integrated laser stabilization, offering a pathway to compact, high-performance photonic oscillators.

Traditional optical isolators rely on magneto-optic materials that are incompatible with CMOS fabrication processes, limiting their integration with silicon photonic circuits. In chapter 4, we present a novel nonmagnetic optical isolator based on acousto-optic interactions in multimode silicon waveguides. By leveraging electrically driven phonon scattering, this approach achieves broadband optical isolation without the need for magneto-optic materials. A multimode waveguide dispersion engineering technique is introduced to extend the isolation bandwidth, resulting in a 2 THz bandwidth with insertion losses as low as -2 dB. The experimental demonstration of this isolator shows significant improvements over existing integrated isolation techniques, making it a viable solution for protecting on-chip lasers and amplifiers from destabilizing back-reflections. This development enhances the robustness of integrated photonic systems, enabling scalable and stable photonic integration.

This dissertation provides a foundation for next-generation silicon photonic technologies by demonstrating novel approaches to laser stabilization and optical isolation. The integration of high-Q reference cavities and broadband nonmagnetic isolators represents a significant step toward fully integrated, high-performance photonic systems

for a wide range of applications.

Chapter 2

Integrated Silicon Nitride Reference Cavity

2.1 Ultra-low-loss Silicon Nitride Waveguide

Silicon nitride waveguides have emerged as a crucial platform for integrated photonics. Due to their low propagation losses, broad transparency window, and compatibility with CMOS processing, Si_3N_4 waveguides are widely employed in areas such as narrow-linewidth lasers, frequency comb generation, precision spectroscopy, and quantum photonics.

Ultra-low loss Si_3N_4 waveguides are essential for realizing high-quality optical resonators that achieve exceptional Q-factors. These high-Q resonators play a pivotal role in stabilizing laser frequencies, which is crucial for applications in optical atomic clocks, communication systems, and quantum information. In addition to frequency stabilization, ultra-low loss Si_3N_4 waveguides enable more efficient nonlinear interactions, such as the generation of optical frequency combs, and can be used in integrated photonic circuits to create narrowband filters with minimal insertion loss. The low propagation loss of these waveguides, achieved through careful control of fabrication

processes, opens up the door for various integrated applications.

Achieving ultra-low propagation loss in Si_3N_4 waveguides requires overcoming several challenges related to the fabrication and material quality of the waveguides. The primary contributors to optical loss in Si_3N_4 waveguides are scattering from surface roughness and absorption due to impurities or imperfections within the material. Scattering loss occurs mostly when imperfections at the waveguide's surface and sidewalls, such as roughness introduced during etching, cause light to couple out of the guided mode and into radiation modes. The degree of this scattering loss is highly sensitive to the roughness of the waveguide's top, bottom, and sidewalls, and its effect is especially pronounced in high-Q resonators, where even minor imperfections can significantly degrade the resonator's performance.

In recent developments, various strategies have been employed to mitigate these losses. High aspect ratio waveguide designs, where the waveguide geometry is optimized to minimize the interaction of the optical mode with the rough sidewalls, have been shown to be effective in reducing scattering losses [5]. Moreover, improvements in lithography and etching techniques have enabled the fabrication of smoother waveguide surfaces, thereby decreasing surface roughness-induced scattering. Additionally, the use of thermal annealing, deposition of a blanket nitride layer, and passivation of absorptive surface states have all contributed to a reduction in material absorption losses, leading to record-low propagation losses for Si_3N_4 waveguides.

Despite these advancements, roughness-induced scattering remains a significant limiting factor in the pursuit of ultra-low loss Si_3N_4 waveguides. Even with state-of-the-art fabrication processes, the formation of roughness during waveguide etching can introduce random defects and irregularities that lead to unwanted scattering and performance degradation. The next section will delve into the mechanisms behind roughness-induced scattering losses in integrated optical waveguides, exploring how sidewall/surface roughness affects the optical mode. By understanding the interplay

between waveguide geometry, surface roughness, and scattering, we can devise new methods to push the limits of waveguide performance, bringing us closer to realizing the full potential of ultra-low loss integrated photonics.

2.2 Roughness-Induced Scattering Loss

In this section, we present a theoretical framework for predicting scattering losses based on the measured roughness data of the waveguide.

2.2.1 Volume current method

The roughness at the boundary of a waveguide can be treated as a perturbation to an ideal smooth waveguide. To calculate the scattered power arising from such perturbations, the volume current method is often employed [28]. Here, we derive and discuss the conventional volume current method in detail.

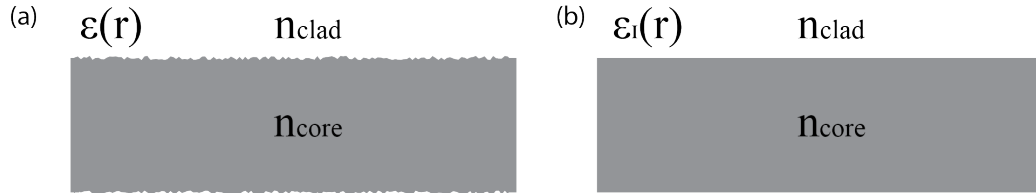


Figure 2.1: Top view of the dielectric distribution for (a) a realistic waveguide and (b) an ideal waveguide.

In a source-free region, Maxwell's equations can be expressed as:

$$\nabla \times (\nabla \times \vec{E}) - \omega^2 \mu \epsilon(\vec{r}) \vec{E} = 0, \quad (2.1)$$

where $\epsilon(\vec{r})$ represents the dielectric distribution of a waveguide, as illustrated in Fig. 2.1(a). This equation can be rewritten as:

$$\nabla \times (\nabla \times \vec{E}) - \omega^2 \mu \epsilon_I(\vec{r}) \vec{E} = \omega^2 \mu (\epsilon(\vec{r}) - \epsilon_I(\vec{r})) \vec{E} = i\omega \mu \vec{J}(\vec{r}), \quad (2.2)$$

where $\varepsilon_I(\vec{r})$ denotes the ideal dielectric distribution of the waveguide, as shown in Fig. 2.1(b).

As evident from the equation above, deviations from the ideal waveguide dielectric profile $\varepsilon_I(\vec{r})$ act as equivalent current sources:

$$\vec{J}(\vec{r}) = -i\omega (\varepsilon(\vec{r}) - \varepsilon_I(\vec{r})) \vec{E} = -i\omega \Delta\varepsilon(\vec{r}) \vec{E}. \quad (2.3)$$

Since the roughness is typically much smaller in scale compared to the waveguide dimensions, $\vec{E}$ can be approximated as the ideal waveguide mode profile $\vec{E}_0$, which satisfies the equation $\nabla \times (\nabla \times \vec{E}_0) - \omega^2 \mu \varepsilon_I(\vec{r}) \vec{E}_0 = 0$. Using these equivalent sources, the radiated power can be calculated with dyadic Green's functions. For cases involving low-index contrast, such as Si_3N_4 and SiO_2 , the Green's function can be approximated by the free-space Green's function. The magnetic vector potential $\vec{A}$ in the far field can be calculated as:

$$\vec{A} = \frac{\mu}{4\pi} \left(\frac{e^{in_{\text{clad}}k_0r}}{r} \right) \int \int \int \vec{J}(\vec{r}') e^{-in_{\text{clad}}k_0\hat{r} \cdot \vec{r}'} dV', \quad (2.4)$$

where k_0 is the free-space wavevector, n_{clad} represents the refractive index of the cladding, and r and $\vec{r}'$ refer to the field and source positions, respectively. The far-field Poynting vector can then be computed as:

$$\vec{S} = \hat{r} \frac{\omega n_{\text{clad}} k_0}{2\mu} |\hat{r} \times \vec{A}|^2, \quad (2.5)$$

with the total radiated power given by:

$$P = \int \int \vec{S} \cdot \hat{r} dA. \quad (2.6)$$

However, one caveat arises when calculating the roughness-induced current source: the field $\vec{E}_0$ is not continuous across the waveguide boundary. Therefore, an important

question arises: should we use the value of $\vec{E}_0(\vec{r})$ as the limit from within or outside the waveguide?

Perturbation theory with shifting material boundary

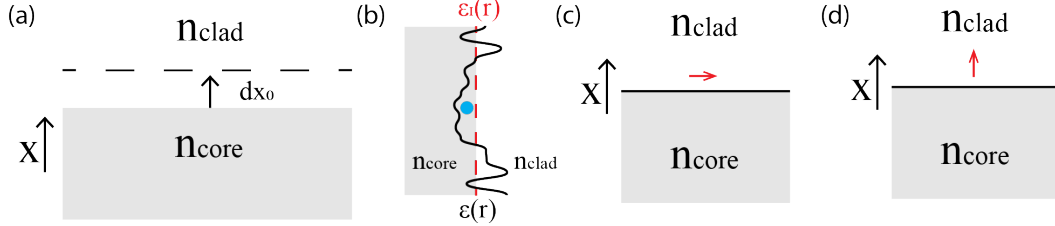


Figure 2.2: (a) Illustration of a perturbation at the dielectric boundary. (b) Comparison between the ideal and perturbed boundaries, highlighting the importance of the permittivity choice at the blue point. (c) Green's function calculation for a dipole near a semi-infinite dielectric boundary with the dipole oriented parallel to the interface. (d) Green's function calculation for a dipole near a semi-infinite dielectric boundary with the dipole oriented perpendicular to the interface.

The issue we mentioned above arises because we introduce a discontinuous transition from n_{clad} to n_{core} during the perturbation of the dielectric boundary, as is shown in Fig. 2.2(a). To address this discontinuity, we can first smooth the dielectric distribution. The original dielectric distribution can be expressed as:

$$\epsilon(x) = \epsilon_{\text{core}} + (\epsilon_{\text{clad}} - \epsilon_{\text{core}})\Theta(x - x_0). \quad (2.7)$$

where $\Theta(x)$ is the Heaviside step function. To smooth the distribution, we can introduce a smoothing function $g_s(x)$, going to a Dirac function $\delta(x)$ as we take the limit of $s \rightarrow 0$.

Boundary continuations of Maxwell equation determine the tangent component of electric field to be contentious whereas the perpendicular component of displacement field to be contentious. For convenience, we smooth the dielectric anisotropically to distinguish tangent and perpendicular components. In this way, we can write the dielectric distribution as:

$$\varepsilon_s(x) = \begin{bmatrix} \tilde{\varepsilon}(x) & 0 & 0 \\ 0 & \bar{\varepsilon}(x) & 0 \\ 0 & 0 & \bar{\varepsilon}(x) \end{bmatrix}, \quad (2.8)$$

with

$$\bar{\varepsilon}(x) = \int g_s(x - x') \varepsilon(x') dx', \quad (2.9)$$

$$\tilde{\varepsilon}(x)^{-1} = \int g_s(x - x') \varepsilon(x')^{-1} dx'. \quad (2.10)$$

Then for a perturbation of dielectric boundary x_0 can be viewed as the perturbation of smoothed relative permittivity value, rather than boundary:

$$\frac{\partial \bar{\varepsilon}}{\partial x_0} = (\varepsilon_{core} - \varepsilon_{clad}) g_s(x - h) \quad (2.11)$$

$$\frac{\partial \tilde{\varepsilon}}{\partial x_0} = -\tilde{\varepsilon}^2 \left(\frac{1}{\varepsilon_{core}} - \frac{1}{\varepsilon_{clad}} \right) g_s(x - h) \quad (2.12)$$

In this way, we can write the volume current as:

$$\vec{J} = -i\omega \left[(\varepsilon_{core} - \varepsilon_{clad}) g_s(x - x_0) \vec{E}_{\parallel} - \left(\frac{1}{\varepsilon_{core}} - \frac{1}{\varepsilon_{clad}} \right) g_s(x - h) \tilde{\varepsilon}(x) \vec{D}_{\perp} \right] \Delta x_0 \quad (2.13)$$

In this form, the tangent electric field $\vec{E}_{\parallel}$ and the perpendicular displacement field $\vec{D}_{\perp}$ are all continuous. Since the perturbation is small enough, we can use the eigenmode of the waveguide, i.e. unperturbed field distribution to calculate the volume current. However, $\tilde{\varepsilon}(x)$ is kind of arbitrary to use the perturbed permittivity or the unperturbed permittivity, while it will make a significant difference proportional to ε_{core} vs ε_{clad} , as is shown in Fig. 2.2(b). To address this contradiction, we need also take the dyadic Green's function into consideration.

For simplicity, we can consider a unit current in the vicinity of a semi-infinite dielectric boundary. First, let's consider the unit current is inside of the cladding. The dyadic Green's function can be easily calculated using image charge method. For a unit current $\hat{J}_{\parallel}$ that is parallel to the interface:

$$\bar{\bar{G}} \cdot \hat{J}_{\parallel} \propto \begin{cases} \frac{2}{\varepsilon_{core} + \varepsilon_{clad}} & x > x_0 \\ \frac{2}{\varepsilon_{core} + \varepsilon_{clad}} & x < x_0 \end{cases} \quad (2.14)$$

For a unit current $\hat{J}_{\perp}$ that is perpendicular to the interface:

$$\bar{\bar{G}} \cdot \hat{J}_{\perp} \propto \begin{cases} \frac{1}{\varepsilon_{clad}} \frac{2\varepsilon_{core}}{\varepsilon_{core} + \varepsilon_{clad}} & x > x_0 \\ \frac{2}{\varepsilon_{core} + \varepsilon_{clad}} & x < x_0 \end{cases} \quad (2.15)$$

where x_0 is the position of the interface. For the case that the unit current is inside of the core, we can just simply swap n_{core} and n_{clad} and domain in the equations above:

$$\bar{\bar{G}} \cdot \hat{J}_{\parallel} \propto \begin{cases} \frac{2}{\varepsilon_{core} + \varepsilon_{clad}} & x > x_0 \\ \frac{2}{\varepsilon_{core} + \varepsilon_{clad}} & x < x_0 \end{cases} \quad (2.16)$$

$$\bar{\bar{G}} \cdot \hat{J}_{\perp} \propto \begin{cases} \frac{2}{\varepsilon_{core} + \varepsilon_{clad}} & x > x_0 \\ \frac{1}{\varepsilon_{core}} \frac{2\varepsilon_{clad}}{\varepsilon_{core} + \varepsilon_{clad}} & x < x_0 \end{cases} \quad (2.17)$$

As we can see, for the unit current that is parallel to the interface, the dyadic Green's function is the same for the two domains, meaning no matter which dielectric the unit current is inside of, the dyadic Green's function is the same. And also note that, for the volume current component that's parallel to the interface $\vec{J}_{\parallel} = -i\omega(\varepsilon_{core} - \varepsilon_{clad})g_s(x - x_0)\vec{E}_{\parallel}\Delta x_0$, we don't have the dilemma of choosing dielectric value.

Now let's look at the perpendicular component. As we mentioned before,

$$\vec{J}_\perp = i\omega \left(\frac{1}{\varepsilon_{core}} - \frac{1}{\varepsilon_{clad}} \right) g_s(x-h) \tilde{\varepsilon}(x) \vec{D}_\perp \Delta x_0$$

is kind of arbitrary to choose $\tilde{\varepsilon}(x)$ as perturbed or ideal while it will lead to completely different value, as is shown in Fig. 2.2(b). However, what really matters is not the current itself, it's the field produced by the current, which is proportional to $\vec{\bar{G}} \cdot \vec{J}_\perp$. Let's use the example shown in Fig. 2.2(b). If we use the ideal dielectric distribution, the current will be:

$$\vec{J}_\perp = i\omega \left(\frac{1}{\varepsilon_{core}} - \frac{1}{\varepsilon_{clad}} \right) g_s(x-h) \varepsilon_{core}(x) \vec{D}_\perp \Delta x_0$$

and the dyadic Green's function will be the case that the current is inside of the core (Eq. 1.17): For $\vec{\bar{G}} \cdot \vec{J}_\perp$ we can see that:

$$\vec{\bar{G}} \cdot \vec{J}_\perp \propto \begin{cases} \frac{2\varepsilon_{core}}{\varepsilon_{core} + \varepsilon_{clad}} & x > x_0 \\ \frac{2\varepsilon_{clad}}{\varepsilon_{core} + \varepsilon_{clad}} & x < x_0 \end{cases} \quad (2.18)$$

For the case of perturbed distribution, the current will be:

$$\vec{J}_\perp = i\omega \left(\frac{1}{\varepsilon_{core}} - \frac{1}{\varepsilon_{clad}} \right) g_s(x-h) \varepsilon_{clad}(x) \vec{D}_\perp \Delta x_0$$

and the dyadic Green's function will be the case that the current is inside of the cladding (Eq. 1.15): For $\vec{\bar{G}} \cdot \vec{J}_\perp$ we can see that:

$$\vec{\bar{G}} \cdot \vec{J}_\perp \propto \begin{cases} \frac{2\varepsilon_{core}}{\varepsilon_{core} + \varepsilon_{clad}} & x > x_0 \\ \frac{2\varepsilon_{clad}}{\varepsilon_{core} + \varepsilon_{clad}} & x < x_0 \end{cases} \quad (2.19)$$

which is consistent with the case of ideal distribution. In conclusion, as long as we

choose the consistent pair of dyadic Green's function and current, the result is the same. To make things convenient, we can make the roughness to be always additive or subtractive compared with the ideal waveguide, i.e the current is always inside of cladding or core.

2.3 Scattering Loss Estimation for Dilute Si₃N₄ Waveguide

Using the volume current method mentioned above, we can estimate the scattering loss induced by surface and sidewall roughness in the dilute Si₃N₄ waveguide [5]. Since sidewall roughness (nanometers) is typically much greater than surface roughness (Angstroms), reducing the core thickness to less than 100 nm helps minimize the optical mode's interaction with the sidewalls. Additionally, the low index contrast between the Si₃N₄ core and cladding further reduces the volume current. To estimate the scattering loss in such a dilute Si₃N₄ waveguide, the free-space dyadic Green's function is a good approximation due to the thin core and low index contrast.

$$\bar{\bar{G}}(r, r') = (1 + \frac{1}{k_0^2} \nabla \nabla \cdot) \frac{1}{4\pi} \frac{e^{ik_0|\vec{r}-\vec{r}'|}}{|\vec{r}-\vec{r}'|} \bar{\bar{I}} \quad (2.20)$$

where k_0 is the free-space wavevector in the cladding. The top surface roughness $f_{top}(x, y)$ can be described by autocorrelation function (using exponential model):

$$R_{top}(u_x, u_y) = \langle f_{top}(x, y) f_{top}(x + u_x, y + u_y) \rangle \approx \sigma_{top}^2 \exp(-\frac{u_x + u_y}{L_{top}}) \quad (2.21)$$

where σ_{top} is the root mean square of the roughness and L_{top} is the correlation length. By incorporating this roughness profile with the volume current method, following the approach in [28], we can get an estimation of scattering loss from the surface roughness profile in dilute Si₃N₄ waveguide.

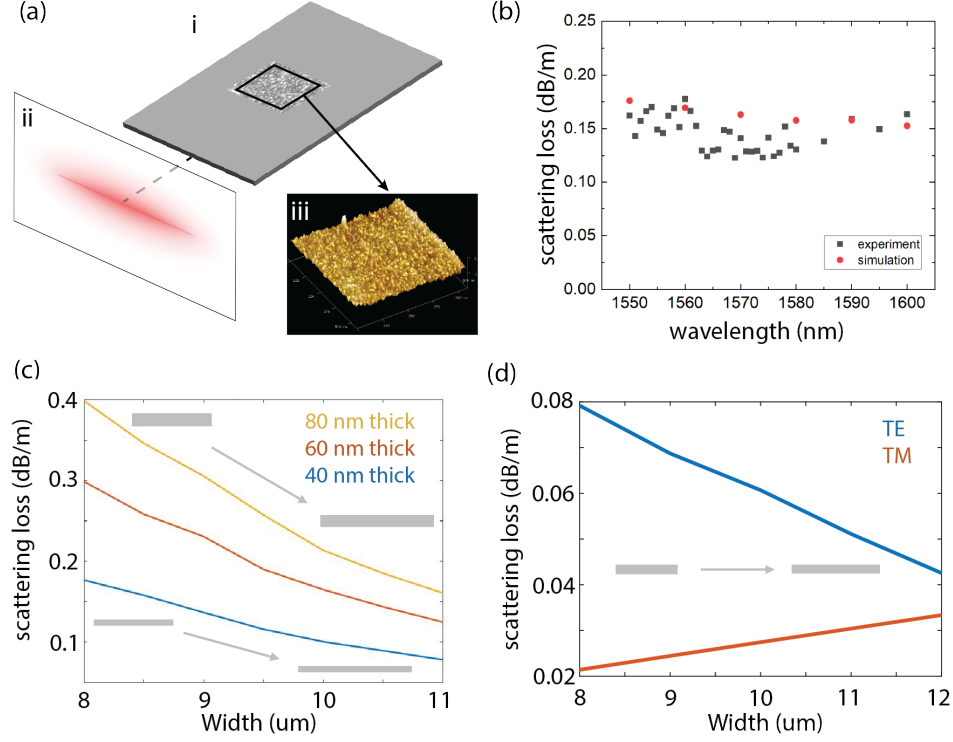


Figure 2.3: (a, i) Schematic of a dilute Silicon Nitride waveguide. (a, ii) Fundamental TE optical mode of the waveguide. (a, iii) Typical surface roughness profile of the Silicon Nitride core. (b) Comparison of the scattering loss estimation to experimental results. (c) Scattering loss prediction for various core geometries. (d) Scattering loss comparison of TE and TM mode.

For the dilute nitride waveguides (Fig.2.3(a)) with an extreme aspect ratio (11μm wide and 40nm thick) [5], simulations reveal that sidewall scattering loss is negligible, while surface scattering loss dominates. Using a typical surface roughness profile of silicon nitride waveguides—0.45 nm RMS roughness and 8.8 nm correlation length, as measured by atomic force microscopy—we simulate scattering loss across various wavelengths. The simulation results align well with experimental data, as illustrated in Fig.2.3(b). This scattering loss simulation framework also serves as a tool for optimizing waveguide structures to minimize scattering loss. For example, Fig.2.3(c) shows simulated scattering loss for various waveguide geometries, indicating that wider and thinner waveguides exhibit lower losses within the explored parameter space. Additionally, Fig.2.3(d) compares scattering losses for different polarization

modes across wavelengths, suggesting the potential for even lower loss when using TM modes instead of TE modes, which we experimentally verified in [29].

2.4 Conclusion

In this chapter, we presented the critical advancements and challenges in achieving ultra-low-loss silicon nitride waveguides. By analyzing roughness-induced scattering losses and employing theoretical framework, volume current method, we have demonstrated tool to estimate how surface and sidewall imperfections impact optical losses. The simulation results also provide a guidance of optimizing waveguide geometries and fabrication processes to minimize scattering losses. Additionally, the comparison between TE and TM modes highlights new opportunities for further reducing propagation loss.

This tool not only works well for dilute Silicon Nitride waveguides, but also applies to other waveguide systems. The only requirement is to use numerical approach to calculate the Green's function in the system [30]. This framework contribute to the development of high-Q resonators for integrated photonics platforms, offering a pathway toward pushing the boundaries of performance in integrated photonic circuits.

Chapter 3

Co-Packaged micro-Fabry Perot Cavities for On-Chip Laser Stabilization

3.1 Introduction

The increasing demand for compact, high-performance laser systems in fields such as telecommunications, quantum technologies, and precision sensing has driven extensive research into integrated photonic platforms. In Chapter 2, we explored the development of high-Q resonators constructed from low-loss silicon nitride waveguides. These resonators have been established as effective on-chip reference cavities, demonstrating remarkable capabilities for laser stabilization through techniques such as Pound-Drever-Hall (PDH) locking [31, 32] and self-injection locking [12]. Despite these advances, silicon nitride-based dielectric resonators face intrinsic challenges, particularly thermorefractive noise, which arises from temperature-induced refractive index fluctuations in the dielectric medium [33]. This noise imposes a fundamental limit on the frequency stability achievable with dielectric resonators, scaling inversely

with the mode size. Consequently, researchers have developed meter-long silicon nitride spiral resonators to mitigate these effects. However, further increasing the mode size is constrained by the limited chip area, and the dense free spectral range (FSR) of such long resonators introduces additional design and operational challenges. Overcoming these challenges requires alternative solutions to achieve ultra-high frequency stability within a compact footprint.

Vacuum gap Fabry-Perot (FP) resonators have many unique properties that make them indispensable as reference cavities. Since they host electromagnetic modes that live almost entirely in vacuum, vacuum gap FP cavities exhibit greatly reduced frequency instabilities relative to dielectric resonators [34–36]. To enable next-generation quantum communications, computation, and timekeeping technologies, it will be necessary to bring the performance advantages of FP resonators to compact, integrated platforms [37] (Fig. 3.1)). As the basis for integrated technologies, new techniques for wafer-scale fabrication of vacuum-gap micro-Fabry-Perot (μ FP) cavities have produced compact reference cavities with ultrahigh Q-factors ($> 10^9$) and excellent frequency stability [13, 38]. However, before we can harness these performance advantages in next-generation integrated photonic systems, new strategies are needed to interface μ FP reference cavities with photonic integrated circuits (PICs).

The most challenging obstacle that prevents the introduction of μ FP into PICs is the born nature of the wide-band and near-unity reflection of FP cavities when off-resonance. Due to the lack of on-chip circulators, such a strong reflection will destabilize and even damage on-chip lasers/amplifiers. To enable the use of μ FP cavities as on-chip frequency references, we require efficient access to the resonator modes and methods to transform the cavity reflective response for various applications. For instance, electrical feedback can be used to actively stabilize a laser to the modes of an FP cavity by detecting the reflection response using Pound-Drever-Hall (PDH)-based feedback control. However, such schemes require the protection of

the on-chip laser from back-reflections and efficient signal redirection for detection. Alternatively, the frequency stability of a μ FP reference cavity can be transferred to a laser oscillator using optical feedback through self-injection locking [39]. This method eliminates complex feedback control while producing noise suppression at higher bandwidths using a simpler system architecture. However, efficient and stable self-injection locking would require new strategies to shape and control the response of the μ FP cavity to enable resonant feedback with integrated laser [14]. Hence, innovative circuit interfaces could facilitate a new generation of ultra-high-performance lasers and oscillators by bringing the performance advantages of vacuum-gap reference cavities to integrated photonics.

In this chapter, we introduce novel reflection transformation circuits (RTCs) for co-integration of vacuum-gap μ FP reference cavities with photonic integrated circuits to enable compact and ultra-high-performance lasers and oscillators. To harness the frequency stability of μ FP reference cavities using PDH-locking schemes, we bring the widely used 'poor man's isolator' scheme from free-space optics into PICs. We demonstrate an isolator-free circuit interface for PDH locking, consisting of a circuit-based two-port interferometer that cancels unwanted back-reflections while mapping the cavity's reflection response to a separate optical port. To harness μ FP cavities for self-injection locking, we demonstrate a new RTC that transforms the reflection response of the μ FP to produce the requisite resonant feedback; using this system, we demonstrate self-injection locking of a waveguide-integrated laser to a vacuum-gap μ FP reference cavity that was fabricated using a wafer-scale process. Self-injection locking of the DFB laser to the μ FP reference cavity, we achieve a single side-band phase noise of -97 dBc/Hz at 10 kHz offset frequency, 5×10^{-13} fractional frequency stability at 10 ms, with a 150 Hz $1/\pi$ integral linewidth and a 35 mHz fundamental linewidth, corresponding to record performance of a self-injection locked laser within a waveguide-integrated system. Looking ahead, these approaches to harnessing refer-

ence cavities within integrated circuits open new avenues for ultra-high performance portable photonic microwave systems [37], integrated optical gyroscopes [40], and enhanced quantum systems [41] with the continued advancement of such μ FP reference cavity technologies.

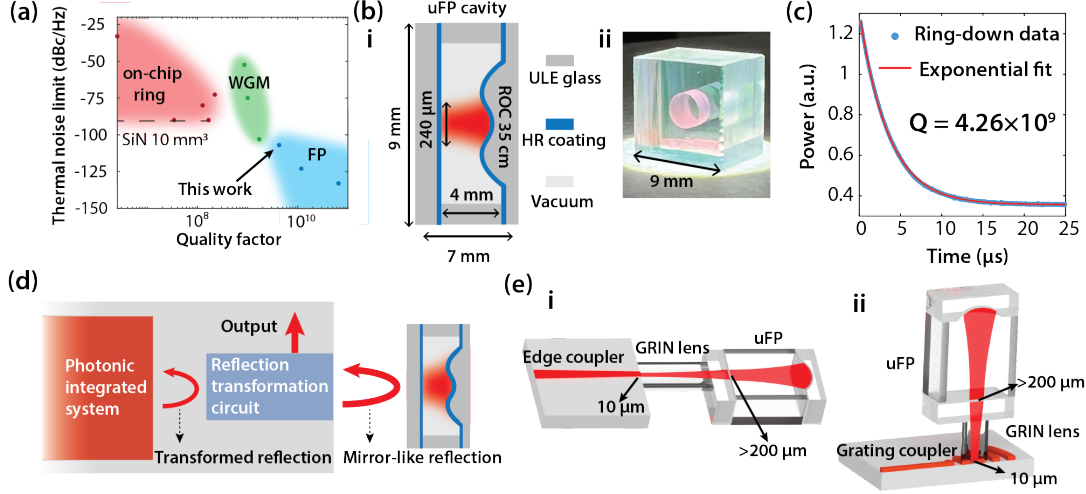


Figure 3.1: **a**, Survey of resonator thermal noise limit at 10 kHz and quality factor for on-chip ring cavity, WGM cavity and vacuum FP cavity. Ring resonator data is from [12,32,42–44], WGM resonator data is from [34,45,46], and vacuum FP resonator data is from this work and [35,47]. Note that our cavity has $> 10^4$ times smaller form factor than [35] and doesn’t require vacuum enclosure. **b**, **i**, Schematic illustration of the in-vacuum bonded μ FP cavity. ROC: radius of curvature. **ii** Photo of μ FP cavity. **c**, Optical ring-down measurement of μ FP cavity yielding quality factor of 4.26 billion. **d** Schematic of the integration architecture. A reflection transformation circuit is required to transform mirror-like reflection to the response desired by application. **e** Schematic of mode-matching approach from PIC to μ FP with **i** end-fire and **ii** grating coupler.

3.2 Co-integration of μ FP Cavities

Integration of high-performance vacuum-gap μ FP reference cavities with waveguide-based PIC platforms presents several challenges. When directly probed, an FP reference cavity will reflect laser light back to its source, which is usually unfavorable for stable laser operation. To solve this issue, we have developed a general on-chip interferometric platform, named reflection transformation circuit (RTC), as is shown

in Fig. 3.1(d). Based on the design, such a platform can either reshape the spectral response of the FP reference cavity reflection that enables on-chip laser self-injection locking, or completely redirect the reflection signal without the need of a circulator for PDH locking. The second challenge is the large difference in size between the modes supported by waveguides and FP reference cavities, where the latter are usually an order of magnitude larger. To overcome this, we demonstrate two complementary strategies with edge and surface couplers to efficiently interface and access millimeter-scale reference cavity modes from micron-scale waveguide modes in two steps. An inverse (Fig. 3.1(e,i)) taper or grating coupler (Fig. 3.1(e,ii)) is first used to expand the waveguide mode to several microns. A micro-optic gradient-index (GRIN) lens is then used to expand and collimate the beam to a mode field diameter of 240 μm , as necessary for efficient mode-matching with the μFP cavity modes.

The vacuum-gap μFP reference cavity used for these studies is shown in Fig. 3.1(b, i-ii), and we use a wafer-scale fabrication process to create such a cavity. A chemically assisted reflow process is used to produce an array of concave micro-mirrors with 35 cm radius of curvature on a super-polished 2-inch glass wafer [36]. Through optical-contact bonding of the concave and flat mirror wafers onto a matched spacer, we construct a μFP cavity array that can be diced into individual cavities (Fig. 3.1(c,ii)). To eliminate unwanted sources of noise and frequency drift, the mirror and spacer are fabricated from ultra-low expansion glass, and the cavity assembly is bonded in a vacuum environment. The resulted plano-concave μFP reference cavity has compact dimensions of $9 \times 9 \times 7$ mm, a small volume less than 0.6 mL and a Gaussian mode field diameter of 240 μm . Using optical ring-down spectroscopy, we measured the optical quality factors of these cavity modes to be 4.26×10^9 (Fig. 3.1(c)), corresponding to a cavity linewidth of 45 kHz and a finesse of 8.25×10^5 . Through independent measurements, such vacuum-gap μFP reference cavities have been shown to produce excellent phase noise and frequency stability [38]. The fabricated μFP cavity with

hermetically sealed vacuum gap offers not only unparalleled noise performance free from the limitations of thermo-refractive noise in dielectric cavities (Fig. 3.1(a)) but also eliminates the need for a vacuum enclosure. Hence, the integration of such high-performance μ FP cavities could bring these benefits to PIC platforms to enable next-generation compact lasers, oscillators, and sensors.

We leverage the integration of a chip-based RTC with μ FP reference cavities to explore two distinct stabilization strategies: electrical feedback and optical feedback. For electrical feedback, we implement an on-chip 'poor man's isolator' scheme within the co-integrated platform to redirect signals and cancel back-reflections effectively. This approach enables the use of μ FP cavities for electrical feedback in a broad range of RF photonic systems. In the optical feedback configuration, we demonstrate direct self-injection locking of an on-chip laser to a μ FP reference cavity. This method utilizes passive optical feedback to significantly enhance laser stability, achieving unparalleled performance in compact integrated photonic platforms. These advancements underline the transformative potential of μ FP cavities in enabling next-generation high-performance lasers and oscillators.

3.3 Reflection Transformation Circuit for PDH Locking

3.3.1 System concept of on-chip poor man's isolator

To effectively enable signal redirection for detection while suppressing back-reflections, we utilize a novel RTC to interface a high-finesse Fabry-Perot resonator for electrical feedback. The operational principle of this photonic circuit interface (Fig. 3.2(b)) draws inspiration from the traditional free-space implementation of the 'poor man's isolator', as depicted in Fig. 3.2(a). In the free-space configuration, p-polarized light

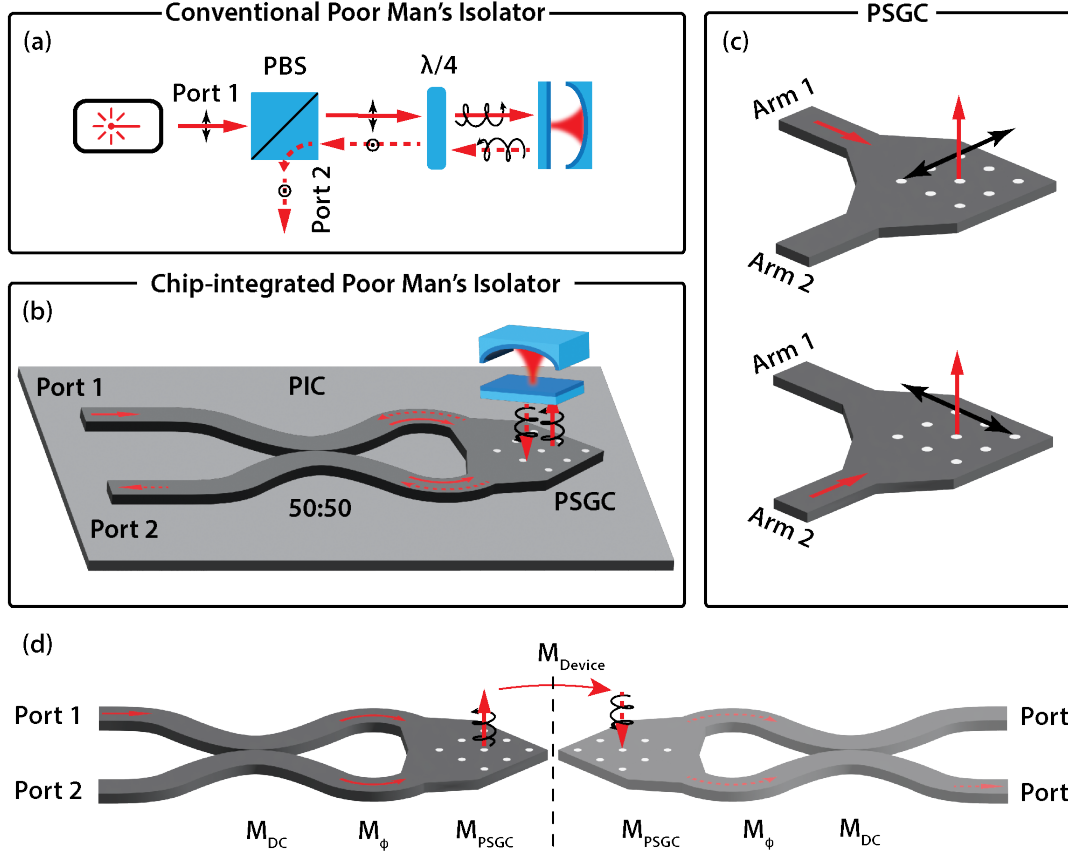


Figure 3.2: **a**, Depiction of the traditional setup of a poor man's isolator mechanism in free space, where PBS denotes a Polarizing Beam Splitter. **b**, Illustration of the circuit interface, which is modeled after the poor man's isolator, that incorporates a the Polarization Splitting Grating Coupler (PSGC). PIC: photonic integrated circuits. **c**, Schematic of PSGC. Light input from Arm 1 will be scattered into p-polarized light in free space and that from Arm 2 will be scattered into s-polarized light in free space. **d** Transfer matrix schematic of the system.

entering Port 1 of the polarizing beam splitter (PBS) passes through and is converted into right-handed circularly polarized light upon traversal through a Quarter Wave Plate (QWP). This light is then reflected by the Fabry-Perot resonator, converting it into left-handed circular polarization. Upon re-entering the QWP, the left-handed circular polarization is transformed into s-polarized light, which is subsequently redirected to exit through Port 2 of the PBS. This setup effectively protects the laser source from back-reflections while ensuring spatial separation of the reflected light, a critical requirement for on-chip PDH locking. This free-space concept has been trans-

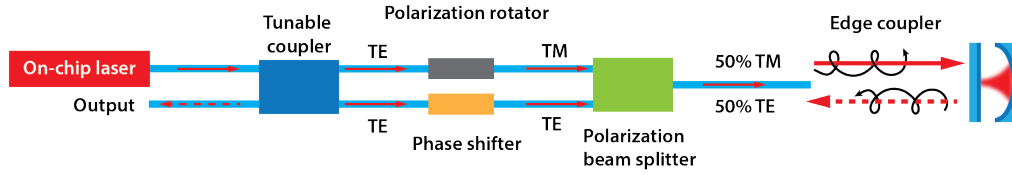


Figure 3.3: Edge-coupler based on-chip poor man's isolator interface circuit.

lated into an integrated photonic circuit that achieves similar functionality using a polarization splitting grating coupler (PSGC) and interferometric design.

The circuit implementation of this poor man's isolator system, seen in Fig. 3.2(b), is comprised of a balanced interferometer that incorporates an optimized polarization-splitting grating coupler (PSGC) device. In close analogy to the free-space polarizing beamsplitter, the PSGC maps orthogonally polarized waves into separate output waveguide arms. The arm 1 (arm 2) of the PSGC couples to p-polarized (s-polarized) free-space beams that are emitted perpendicular to the grating coupler (Fig. 3.2(c)). Light entering Port 1 of this interferometer is split between two waveguides by a 50/50 directional coupler before coupling to the PSGC. Since the directional coupler induces a $\pi/2$ phase difference between the two waveguides (each having identical path lengths), right-handed circularly polarized light is emitted from the PSGC. Upon reflection from the Fabry-Perot resonator (which is assumed to have a polarization-independent reflection response), the incident right-handed circular polarization is converted into a left-handed circular polarized wave before entering the PSGC for a second time. This left-handed circular polarized wave is projected into orthogonal linear polarizations by the PSGC, meaning that the waves exiting the PSGC now have a $-\pi/2$ phase difference. This phase difference causes these two reflected waves to combine within the 50/50 directional coupler, such that all of the reflected light exits Port 2 of the interferometer. Hence, this circuit interface yields spatial separation of the incident and reflected waves (as required for on-chip PDH locking) while also

producing interferometric cancellation of unwanted back-reflections from Port 1.

Next, we use the transfer matrix formalism to analyze the response of this circuit interface. By mapping the bi-directional two-port system of Fig. 3.2(b) onto an equivalent unidirectional two-port system of Fig. 3.2(d), we can readily use 2×2 transfer matrices to analyze the system response. Denoting the incident waves as $\begin{bmatrix} a_1 \\ a_2 \end{bmatrix}$ and the out-going waves as $\begin{bmatrix} b_1 \\ b_2 \end{bmatrix}$, we can find the system response by propagating the input waves through transfer matrices associated with each component of the system, as follows:

$$\begin{bmatrix} b_1 \\ b_2 \end{bmatrix} = \mathbf{M}_{\text{DC}} \mathbf{M}_\phi \mathbf{M}_{\text{PSGC}} \mathbf{M}_{\text{Device}} \mathbf{M}_{\text{PSGC}} \mathbf{M}_\phi \mathbf{M}_{\text{DC}} \begin{bmatrix} a_1 \\ a_2 \end{bmatrix}. \quad (3.1)$$

Here, $\mathbf{M}_{\text{DC}} = \begin{bmatrix} r & -i\mu \\ -i\mu & r \end{bmatrix}$, $\mathbf{M}_\phi = \begin{bmatrix} 1 & 0 \\ 0 & e^{i\phi} \end{bmatrix}$, $\mathbf{M}_{\text{PSGC}} = \begin{bmatrix} \alpha & 0 \\ 0 & \alpha \end{bmatrix}$, and $\mathbf{M}_{\text{Device}} = \begin{bmatrix} \tilde{R}_1 & 0 \\ 0 & \tilde{R}_2 \end{bmatrix}$ are transfer matrices that describe the response of the directional coupler, differential phase delay, PSGC and reflective device, respectively. Here, μ and $r = \sqrt{1 - \mu^2}$ are splitting coefficient of the directional coupler, ϕ is the phase imbalance between two waveguide segments, α is the coupling coefficient for PSGC, and $\tilde{R}_1, \tilde{R}_2$ are complex reflection coefficients for the external device (e.g., Fabry-Perot resonator) for s- and p-polarized light.

Using this model, we make some basic observations about this reflection cancellation scheme. Assuming that light is only injected into Port 1 (i.e., $a_2 = 0$), we find reflected wave amplitudes $b_1 = a_1 \alpha^2 (\tilde{R}_1 r^2 - \tilde{R}_2 \mu^2 e^{2i\phi})$ and $b_2 = -a_1 \alpha^2 r \mu (\tilde{R}_1 + \tilde{R}_2 e^{2i\phi})$, from the above transfer matrix model. From these expressions, we see that it is possible to null b_1 even in cases when $\tilde{R}_1 \neq \tilde{R}_2$. However, through experimental studies, we are interested in the case when the reflection response is identical for both polarizations. Assuming that $\tilde{R}_1 = \tilde{R}_2$ in the case of 50% power splitting ratio ($r = \mu = 1/\sqrt{2}$), we see that it is possible to null b_1 while directing all of the output light b_2 in the case when $\phi = 0$. Note, however, that the backreflection suppression ratio, $|b_2/b_1|^2 = \frac{4\mu^2(1-\mu^2)}{|1-2\mu^2|^2}$, is sensitive to splitting coefficient μ . For example, if μ devi-

ates from $1/\sqrt{2}$ by 14% due to fabrication errors, the backreflection suppression ratio will decrease to 10 dB from perfect cancellation. Hence, tight control of the power splitting ratio is necessary to obtain a high degree of back-reflection cancellation.

With the same principle, the poor man’s isolator can also be implemented using edge coupler-based circuits, as shown in Fig. 3.3. In this design, the circuit leverages the TE and TM modes in optical waveguides to map the two orthogonal polarizations in free space. In one arm, an integrated polarization rotator converts the TE mode into a TM mode. The TE and TM modes are then recombined using an on-chip polarization beam splitter, creating an output with 50% TE mode and 50% TM mode at the edge coupler. This configuration effectively generates circularly polarized light, enabling a compact and robust approach to optical isolation within photonic integrated circuits.

3.3.2 Inverse-design optimized polarization splitting grating coupler

To realize this circuit interface, we fabricate a silicon photonic circuit from a Silicon-On-Insulator (SOI) wafer having a 250 nm thick silicon layer and a 3 micron silica under cladding. We use e-beam lithography and a reactive-ion etch (etch depth of 80 nm) to define both waveguide and grating structures, as seen in Fig.2(a). To accurately define the desired structure during e-beam exposure, we implemented proximity effect correction [48], and through dose tests using hydrogen silsesquioxane (HSQ), an e-beam dose of $1050 \mu\text{C}/\text{cm}^2$ yielded the optimal performance. For details of the fabrication process, see [49]. The PSGC device is a key component of this system, as the efficiency of circuit interfaces hinges on the performance of this grating coupler.

To optimize the efficiency of the PSGC, we employ an inverse design algorithm [50] using LUMOPT package with Ansys Lumerical software. Through the design of the grating coupler, the parameter space of numerical optimization is reduced by impos-

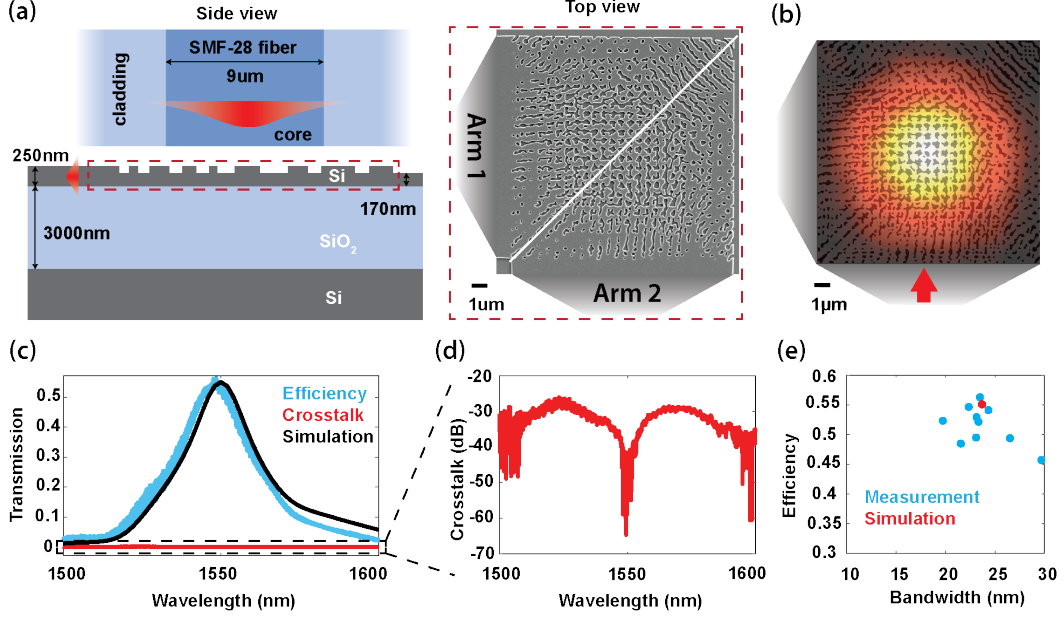


Figure 3.4: **a**, Side view: The layered structure of the PSGC is presented, with a single mode optical fiber positioned perpendicularly to the PSGC. Top view: SEM image of the PSGC. Pattern is symmetrical over the white line. **b**, FDTD simulation showcases the scattered field 3 μm above the Si device layer, prompted by an incident waveguide mode from bottom port. **c**, Empirical measurement of the coupling efficiency and crosstalk of the PSGC. Efficiency is characterized as the transmission from Arm 1 to the fiber, while crosstalk is defined as the transfer from Arm 1 to Arm 2. **d**, Zoom-in of the crosstalk data presented in (c). **e**, Scatter plot showing the performance among ten distinct PSGCs fabricated on the same chip.

ing a mirror symmetry plane, indicated by the diagonal white line in Fig. 3.4(b). We modified LUMOPT package in order to impose such mirror symmetry. This symmetry plane ensures that the grating coupler produces the same scattering response when excited from either input port. The PSGC structure (seen in Fig. 3.4(a-b)) is designed for vertical coupling to SMF-28, with a Gaussian beam waist radius of $5.2\mu m$. Fig. 3.4(b) shows the finite-difference time-domain (FDTD) simulation of the Gaussian field profile produced by the PSGC device layer when illuminated by a TE-like guided optical mode.

To evaluate the optical performance of the fabricated PSGC structures, a cleaved optical fiber (SMF-28) is vertically aligned with the PSGC to enable fiber-to-chip coupling efficiency and cross-talk measurements. The optical measurements seen in

Fig. 3.4(c), reveal fiber-to-chip coupling efficiency of 55% (2.6 dB) at a wavelength of 1550 nm, over a 23 nm bandwidth. The crosstalk, defined as the direct coupling from Arm 1 to Arm 2, displays a contrast exceeding 60 dB at 1550 nm. Fig. 3.4(e) shows the measured peak efficiency and 3 dB bandwidth of 10 gratings fabricated on the same chip, revealing high efficiency (>50%) and broad operational bandwidth (>20 nm) for the majority of fabricated gratings. Deposition of a mirror underneath PSGC will further help to increase to efficiency to nearly 100%, avoiding light scattered into substrate [51].

3.3.3 Experimental results

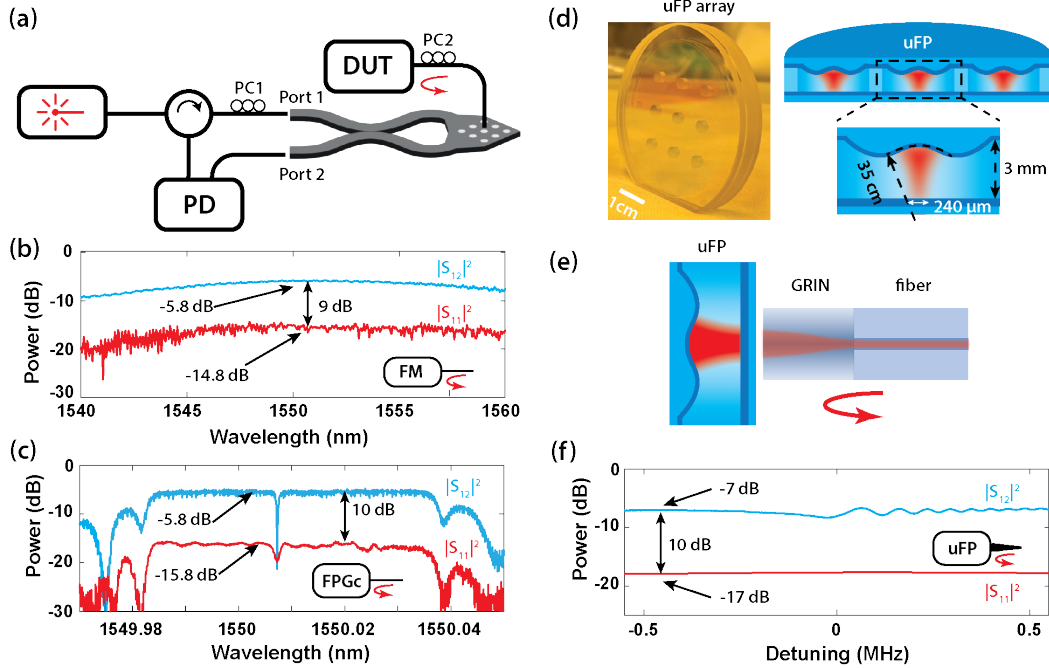


Figure 3.5: **a**, Measurement setup for the characterization of the reflection cancellation circuit. DUT: device under test, PD: photo detector. **b-c**, Measurement results with fiber mirror (FM) and fiber Bragg grating cavity (FBGC) as DUT. **d**, μ FP array. **e**, Fiber GRIN lens collimator assembled to match μ FP cavity mode. **f**, Measurement results with fiber-GRIN- μ FP as DUT.

To validate the proposed RTC and its ability to enable reflection cancellation and signal redirection, we conducted a series of experiments interfacing the RTC with var-

ious reflective devices. Fig. 3.5(a) shows a schematic of the experimental apparatus that is used to assess the performance of the RTC. The fiber circulator facilitates the measurement of system reflection, enabling the characterization of the back-reflection suppression ratio. Through these studies, we use a segment of SMF-28 fiber as a fiber umbilical to flexibly interface the grating coupler to different reflective devices. This fiber umbilical contains a polarization controller (PC2) that compensates for any polarization distortion occurring within the fiber. The state of the polarization controller is chosen to ensure that the fiber umbilical produces a Jones matrix of identity (i.e., doesn't alter the polarization state). Hence, the fiber umbilical is a convenient way to couple to different devices under test (DUT) to evaluate the performance of this circuit interface in various scenarios.

We begin by using a commercial fiber mirror (FM) as a reference device (Thorlabs P5-SMF28ER-P01-1) to examine the performance of the RTC. The reflection and transmission response produced by the fiber mirror are shown in Fig. 3.5(b). Power transmission from Port 1 to Port 2, denoted as $|S_{12}|^2$, reveals a signal attenuation of -5.8 dB. Since the light exiting Port 2 passes through the grating coupler twice, with 2.6 dB of loss per pass, and the fiber mirror has a typical reflection loss of 0.6 dB, the total transmission loss of 5.8 dB is in good agreement with the anticipated transmission loss of 5.8 dB. Measurements of the reflection response from Port 1, denoted as $|S_{11}|^2$, reveal a back-reflection of -14.8 dB from a fiber mirror with near-unity back-reflection efficiency. Since the reflection efficiency, $|S_{11}|^2$, is 9 dB lower than the transmission efficiency, these measurements reveal a back-reflection suppression ratio of 9 dB, validating the operating principle of this reflection cancellation circuit.

Replacing the fiber mirror with a fiber Bragg grating cavity (FBGC), we next examine how the reflection response of this Fabry-Pérot resonator is imprinted on the response of this multi-port system. This FBGC is a commercial fiber Bragg defect cavity (Teraxion C56514-0005), that consists of two fiber-Bragg grating mirrors

that support a standing-wave mode. The cavity produces a single resonance at a 1550 nm center wavelength having a linewidth of 85 MHz and Q-factor of 2.3 million. Fig. 3.5(c) displays the reflection, $|S_{11}|^2$, and transmission, $|S_{12}|^2$, measurements produced when this FBGC is coupled to the circuit. We see that the reflection response of the resonator is clearly imprinted on the transmission response as required for PDH-locking of a laser to the cavity mode. A transmission efficiency of -5.8 dB is achieved (from input to output port), while the cancellation circuit reduces back-reflection to -15.8 dB, corresponding to a back-reflection a back-reflection suppression ratio of 10 dB.

Next, we couple this circuit to a compact, ultra-high finesse micro-Fabry-Pérot (μ FP) resonator that could serve as a chip-integrated frequency reference for future compact optical clock technologies. This μ FP cavity (see Fig. 3.5(d)) is one of an array of air-gap resonators that was created using a wafer-scale reflow fabrication process described in Ref. [13]. Through independent cavity ring-down spectroscopy measurements, the μ FP of interest was found to have an optical linewidth of 60 kHz and an optical Q-factor of 3.2 billion. This resonator was designed to support a relatively large Gaussian-beam diameter of 240 μ m to reduce the impact of thermal noise imparted by the mirror coating on the cavity mode. As a demonstration of the potential for integration of a μ FP cavity with the chip, the 10.4-micron fiber mode was efficiently coupled to the 240 micron resonator mode using a compact (1mm $\times$ 0.5mm) GRIN lens.

During transmission measurements (see Fig. 3.5(f)), the laser sweeps through the cavity resonance over a time interval that is much shorter than cavity ringdown time, transforming the resonant response of the cavity into a ringdown response. This type of ringdown response is due to the interference between energy transiently stored in the cavity and the transmitted laser light [52–55]. The transmission response, $|S_{12}|^2$, reveals the ringing cavity response atop a transmission response showing a 7

dB average transmission loss, consistent with a small amount of excess loss produced by fiber-GRIN coupling to the resonator. Back-reflection, $|S_{11}|^2$, of -17 dB seen in Fig. 3.5(f) corresponds to a back-reflection suppression ratio of 10 dB, demonstrating the suppression of unwanted back-reflections using such reflection cancellation circuits. As discussed in Section IV, a higher level of integration could be achieved by eliminating the fiber umbilical, and directly bonding the GRIN-cavity assembly to the grating coupler.

Building on these results, higher grating coupler efficiencies and lower optical losses could be realized by creating a highly reflecting back-plane underneath PSGC. Moreover, further improvement of the back-reflection suppression ratio could be achieved by using tunable directional couplers and phase shifters, which are readily available in Si photonics foundry processes [56], to tune the parameters of the circuit to obtain a much higher (>30 dB) degree of reflection cancellation.

3.3.4 Co-packaged μ FP cavity with poor man's isolator interface through PSGC

Building on the principles of RTC for signal redirection and reflection cancellation, we now examine the co-packaging of μ FP cavities with PICs (Fig 3.6(a)). Efficient coupling into the μ FP cavity is achieved using a GRIN lens, which expands the beam emitted by the PSGC. The PSGC is designed to produce vertical emission of a 10.4 μ m mode field diameter. This beam is expanded to produce a collimated 240 μ m mode field diameter at the input μ FP cavity using a GRIN lens with an effective focal length of 0.94 mm. The GRIN lens and μ FP cavity are then co-integrated with PIC using a custom spacer that provides mechanical stability for the bonded assembly, as is shown in Fig. 3.6(b). Using this integration method, light can be coupled from the silicon circuit into the μ FP cavity with an efficiency of -2.8 dB. A fiber array was bonded the chip-integrated μ FP module to interrogate the response of this multi-port

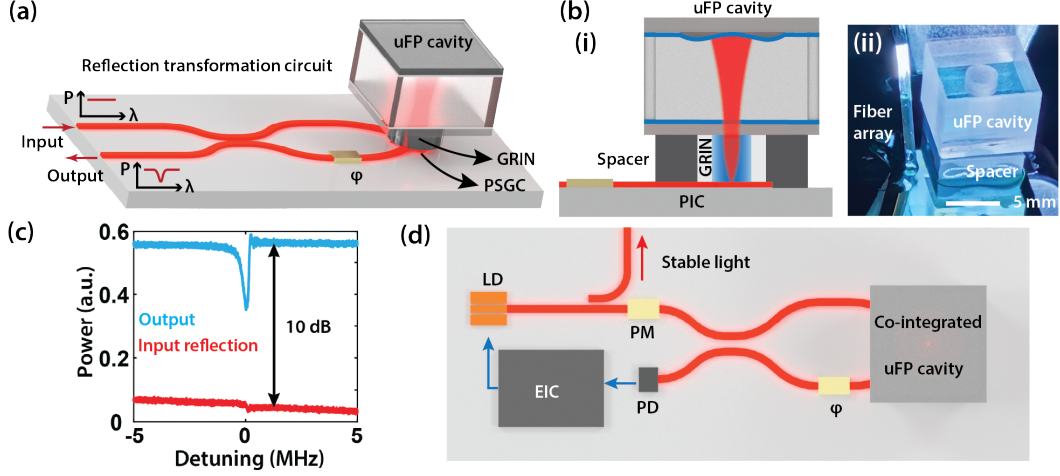


Figure 3.6: **a**, Schematic of the approach for co-integrating μ FP cavity with signal redirection and back-reflection isolation. **b**, **i** Cross section of co-integrated μ FP cavity. **ii** Optical picture for co-integrated module. **c**, Experimental data validating signal redirection and characterizing back-reflection isolation ratio. **d**, Outlook for on-chip PDH locking system. PM: phase modulator. PD: photo detector. LD: laser diode. EIC: electric integrated circuits.

photonic interface.

Fig. 3.6(c) shows the optical reflection from Port A and the transmission to Port B as a function of laser detuning from the cavity resonance. As seen in Fig. 3.6(c), this circuit maps the μ FP cavity reflection response to output port while simultaneously suppressing back-reflections from input port. The relative magnitude of the transmission and reflection reveal 10 dB back-reflection suppression, to aid in protecting an on-chip laser from unwanted feedback. Such a fully integrated module is readily suitable for an on-chip PDH locking system (Fig. 3.6(d)) as proposed in [37]. Together with many other applications such as photonic RF filtering and microwave signal generators, the co-integration of vacuum-gap μ FP cavities paves the way for a next-generation of compact, high-performance RF photonic systems.

3.4 Reflection Transformation Circuit for Self-injection Locking

3.4.1 System concept

Self-injection locking (SIL) utilizes optical feedback to lock a laser frequency's to an external cavity, causing the laser to inherit the noise characteristics of the cavity. Recent studies have used silicon nitride ring resonators to self-injection lock chip-scale semiconductor lasers, greatly improving their phase noise characteristics [11,42,47,57]. These demonstrations have relied on Rayleigh back-scattering to produce resonant feedback necessary for self-injection locking. However, thermo-refractive noise and scattering losses limit the achievable quality factors and noise characteristics of such waveguide-based resonators [12,29]. Alternatively, if such dielectric ring resonators can be replaced by vacuum-gap μ FP reference cavities, thermorefractive noise could be practically eliminated, offering a path to lasers with greatly improved phase noise and frequency stability. However, self-injection locking of integrated lasers will require new strategies to interface such μ FP cavities with photonic integrated circuits to produce resonant feedback.

Established methods for producing resonant feedback have relied on circulators to redirect the transmission response of the FP cavity into the laser [58]. However, this strategy would be challenging to implement using photonic integrated circuits, as isolators and circulators are not widely available; moreover, it would be difficult to access both ports of the μ FP cavity using a single chip. Alternatively, resonant feedback can be produced by partial excitation of higher order modes using a tilted cavity configuration; however, this approach requires a more complex alignment procedure, and it results in excess coupling losses that unnecessarily reduce the efficiency of laser feedback [59–62].

To address these challenges, we demonstrate a new circuit interface that trans-

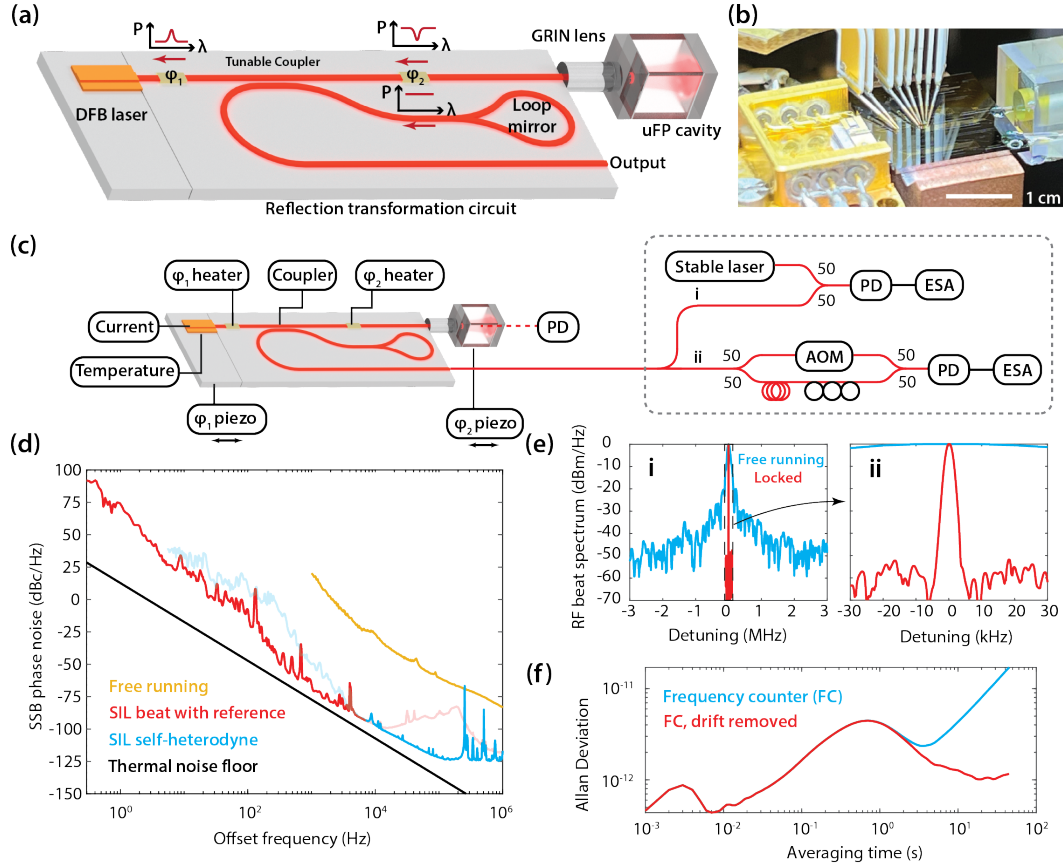


Figure 3.7: **a**, Schematic illustration of on-chip circuit design enabling self-injection locking to co-integrated μ FP cavity. **b**, Photograph of DFB laser self-injection locked to μ FP cavity through interface chip. **c**, Schematic of experimental setup for noise measurement. Black lines indicates electrical control, and red lines indicates optical circuits. PD: photo detector. ESA: electrical spectrum analyzer. AOM: acousto-optical modulator. The output of SIL laser either beats with another stable laser or performs self-heterodyne with 800 m delay line for the ESA to measure phase noise. **d**, Phase noise spectrum for SIL and free running laser. Free running laser phase noise is measured using self-heterodyne. The faded curve indicates the noise is limited by the measurement noise floor. **e**, RF beat note spectrum of the SIL laser and a stable laser reveals linewidth compression when locked. **ii** is the zoom in for **i**. Free running beat note is taken with 91 kHz RBW, and locked beat note is taken with 3 kHz RBW. **f**, Allan deviation of the SIL laser measured from frequency counter with and without linear drift removal.

forms the reflection response of the μ FP cavity to produce resonant feedback necessary for self-injection locking. In reflection, the cavity modes of an FP resonator appear as anti-resonances (dips) atop a high reflectivity background; such anti-resonances are unsuitable for self-injection locking. These antiresonances result from interfer-

ence between light that is resonantly scattered by the cavity modes and a broadband background reflection produced by the mirrors. To utilize the μ FP for self-injection locking, it is crucial to eliminate the high reflectivity background, leaving only resonant scattering (or cavity leakage field) to produce feedback; this is because broadband mirror reflections can produce parasitic resonances that destabilize the laser [14]. If the wideband mirror reflection can be eliminated, leaving only resonant backscatter from the cavity modes, efficient resonant feedback can be achieved. In this case, the sudden phase shift produced by resonant backscattering about the cavity resonance provides the optimal feedback mechanism for self-injection locking.

The circuit interface that we use to produce resonant feedback into a waveguide-based DFB laser is seen in Fig. 3.7(a). An on-chip directional coupler divides the incoming light, directing a portion of light into an on-chip loop-mirror that introduces an additional broadband mirror reflection. The remaining light travels to the chip facet, where it is collimated by a GRIN lens to match the mode of the co-integrated FP cavity. These two reflections interfere upon recombination at the directional coupler. With the correct phase ϕ_2 , the loop-mirror reflection destructively interferes with the FP cavity reflection, cancelling the broadband mirror reflections and preserving only the resonant backscattered light. This results in a resonant (Lorentzian) feedback spectrum, facilitating self-injection locking by adjusting the feedback phase ϕ_1 .

For the choice of GRIN lens, it is crucial to expand and collimate the beam of light emitted from the waveguide end-facet to match the μ FP cavity mode. To accomplish this mode transformation using our circuit interface, we use an inverse taper at the output facet of a Si_3N_4 waveguide, producing a mode diameter of approximately $3\text{ }\mu\text{m}$. A GRIN lens with an effective focal length of 0.46 mm is then used to expand this $3\text{ }\mu\text{m}$ to $240\text{ }\mu\text{m}$, matching the mode field diameter of the μ FP cavity, permitting coupling losses as low as 1.5 dB . Hence, this simple and compact photonic interfacing technique, shown in Fig. 3.7(b), offers a practical means of co-integrating the μ FP

cavity with PICs, to enable high performance oscillators and lasers for next-generation photonic systems.

The circuit interface of Fig. 3.7(c) is used to control the optical feedback parameters of this system, as necessary to obtain a stable self-injection locking. The interface circuit includes integrated heaters that allow for adjustment of the coupler splitting ratio, the feedback phase ϕ_1 , and the interferometer phase ϕ_2 . Piezo actuators on both the DFB laser and the μ FP cavity are also used for fine-tuning ϕ_1 and ϕ_2 by modulating their separation from the photonic circuit. When ϕ_1 and ϕ_2 are adjusted to produce a self-injection locked state, the frequency of the laser emission becomes resonant with the μ FP cavity, producing a dramatic increase in the light intensity within the μ FP cavity. Light transmitted from the μ FP cavity is collected through free space and monitored with a photodetector to observe the onset of the self-injection locking process. The output from the SIL laser is then coupled into the fiber-based apparatus of Fig. 3.7(c,i-ii) for phase noise measurements.

3.4.2 SIL laser with co-packaged μ FP cavity

Two complementary techniques are used to quantify the phase noise of the self-injection-locked laser. To measure the phase noise at low offset frequencies, we analyze the beat note between the integrated SIL laser and another frequency-stabilized reference laser using a phase noise analyzer, as shown in Fig. 3.7(c,i). The stable laser is comprised of a fiber laser that is PDH-locked to a 10 cm long table-top FP reference cavity housed in a vacuum chamber that is known to have low phase noise at low offset frequency (0 dBc/Hz at 1 Hz offset frequency). However, since the PDH feedback loop introduces excess noise, seen as a servo bump at offset frequencies above 10 kHz, a complementary measurement method is required to analyze the phase noise of SIL laser at higher offset frequencies. Here we use the fiber interferometer with 800 m delay seen in Fig. 3.7(c,ii) to perform self-heterodyne measurements. This com-

plementary approach is ideal for phase noise measurements at high offset frequencies since the fiber interferometer is essentially free of technical noise for frequencies above 5 kHz. Combining these two techniques, we obtain a complete picture of phase noise spectrum of the SIL laser from 1Hz to 1MHz, as seen in Fig. 3.7(d).

The measurements of Fig. 3.7(d) reveal a dramatic reduction in the phase noise of the DFB laser when it is injection locked to the μ FP cavity. Self-injection locking is seen to reduce the phase noise of the DFB laser by eight, seven, and six orders of magnitude at offset frequencies of 1 kHz, 10 kHz, and 100 kHz, corresponding to phase noise levels of -65 dBc/Hz, -97 dBc/Hz, and -122 dBc/Hz, respectively. Further analysis of the phase noise spectrum reveals a fundamental linewidth of 35 mHz and a $1/\pi$ integral linewidth of 150 Hz. The reduction in laser linewidth produced by self-injection locking is further demonstrated in the direct RF beat note spectrum between the SIL laser and a stable reference laser, as shown in Fig. 3.7(e). It is important to note that the linewidth in Fig. 3.7(e,ii) is constrained by the resolution bandwidth (RBW) of the electronic spectrum analyzer. Note that these measurements were performed with a relatively large resolution bandwidth (>3 kHz) to ensure a stable beat note spectrum within the acquisition time. The Allan deviation, which is indicative of fractional frequency stability for different averaging times, is seen in Fig. 3.7(f), revealing 5×10^{-13} at 10 ms averaging time. After removing linear drift of 109 Hz/s, the system exhibits a fractional frequency stability better than 5×10^{-12} from 1 ms to 50 s averaging times.

Note that these measurements correspond to record-level phase noise among state-of-the-art chip-integrated systems [12] and were obtained without tight environmental controls. The measurements described above were obtained without active temperature stabilization of the cavity, illustrating the temperature insensitivity and robustness of such in-vacuum bonded μ FP reference cavities. Deviations from the thermal noise limit at low offset frequencies are attributed to mechanical noise from the trans-

lation stages used to position the cavity and DFB laser. Hence, the low offset phase noise and frequency drift are likely to greatly improve when this system is fully co-integrated.

Dynamic of μ FP - SIL laser

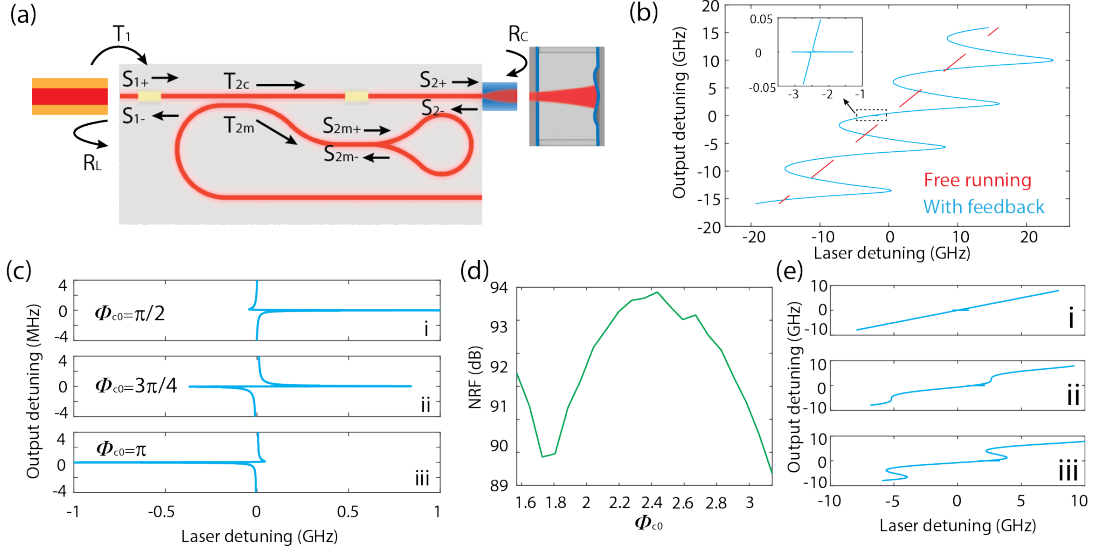


Figure 3.8: **a**, Schematic of the system. **b**, Response curve for a DFB laser directly coupled to a μ FP cavity. **c**, Response curve for a DFB laser directly SIL to a μ FP cavity with perfect mirror reflection removal, shown for various feedback phases. **d**, Noise reduction factor for a DFB laser directly SIL to a μ FP cavity with perfect mirror reflection removal as a function of the feedback phase. **e**, Response curve for a DFB laser directly SIL to a μ FP cavity with imperfect mirror reflection removal. Cases (i–iii) correspond to $\phi_{m0} - 2\phi_{c0} = 0, -\pi/20, -\pi/10$.

In this subsection, we describe the dynamics of self-injection locking to the μ FP cavity. The laser and cavity fields are denoted as A_L and A_C . Due to the multiple reflection interfaces (potentially forming etalons), we use a scattering matrix approach to describe the response of the interface circuit, as shown in Fig. 3.8(a). The equations of motion can be written as [42]:

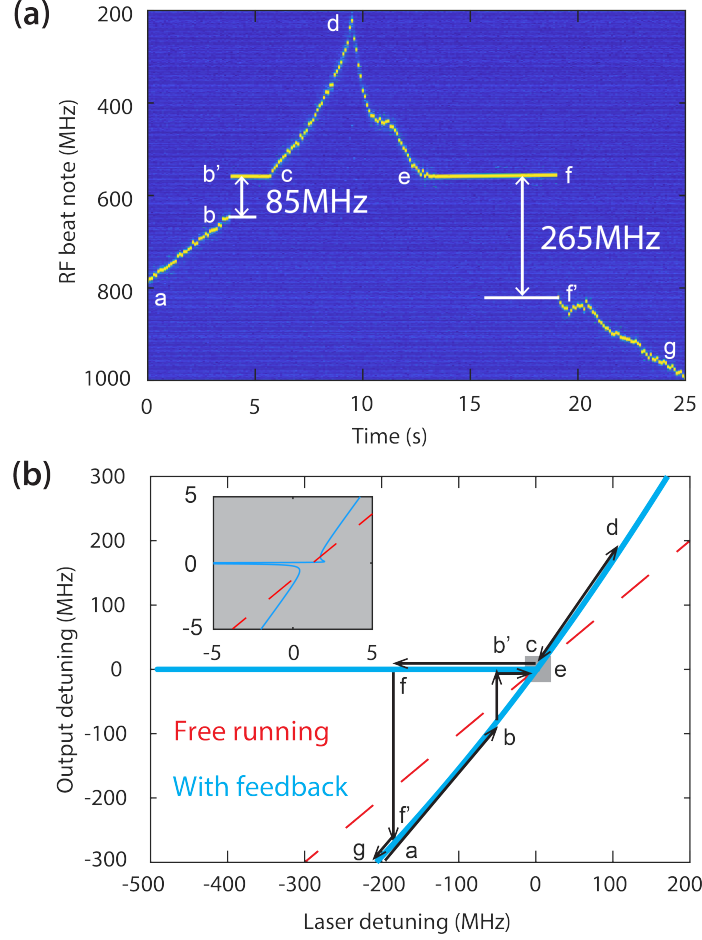


Figure 3.9: **a**, RF beat note spectrum of the SIL laser and a stable laser with a sweeping current on the DFB laser yields a 265 MHz locking range. **b**, Theoretical locking curve. The SIL laser entering and exiting the locked state are due to relaxation oscillation of the DFB laser. Off-resonance curve deviates from free running due to the residual broadband mirror reflection.

$$\frac{dA_L}{dt} = i(\delta\omega_L - \delta\omega)A_L - \frac{\gamma_L}{2}A_L + \frac{g_L}{2}(1 + i\alpha)A_L + i\sqrt{\kappa_L T_1}S_{1-}, \quad (3.2)$$

$$S_{1+} = \sqrt{T_1 R_L}S_{1-} + i\sqrt{\kappa_L T_1}A_L, \quad (3.3)$$

$$S_{2+} = S_{1+} \sqrt{T_{2c}} e^{i\phi_c}, \quad (3.4)$$

$$S_{2-} = \sqrt{R_C}S_{2+} + i\sqrt{\kappa_C}A_C, \quad (3.5)$$

$$\frac{dA_C}{dt} = -i\delta\omega A_C - \frac{\gamma_C}{2}A_C + i\sqrt{\kappa_C}S_{2+}, \quad (3.6)$$

$$S_{1-} = \sqrt{T_{2c}}S_{2-} e^{i\phi_c} + i\sqrt{T_{2m}}S_{2m-} e^{i\phi_{m2}}, \quad (3.7)$$

$$S_{2m-} = S_{2m+} e^{i\phi_{m1}}, \quad (3.8)$$

$$S_{2m+} = iS_{1+} \sqrt{T_{2m}} e^{i\phi_{m2}}. \quad (3.9)$$

Here, T_1 is the energy coupling efficiency from the laser to the photonic integrated circuit (PIC). The quantities T_{2c} and T_{2m} are energy coupling coefficients of the on-chip directional coupler to the cavity and the loop-mirror side, respectively, with the constraint $T_{2c} + T_{2m} = 1$. The parameters R_L and R_C are the reflectivities of the laser mirror and the cavity mirror. The term $\delta\omega_L$ is the detuning of the laser cavity relative to the μ FP cavity, and $\delta\omega$ is the detuning of the output light relative to the μ FP cavity. The parameter g_L is the intensity-dependent gain of the laser, and α is the amplitude-phase coupling coefficient. The decay rates of the laser cavity and the μ FP cavity are denoted by γ_L and γ_C , respectively, while κ_L and κ_C are their respective external coupling rates. ϕ_c is the accumulated phase from the input of the PIC to the μ FP cavity side. Finally, ϕ_{m1} and ϕ_{m2} are the phases accumulated in the loop-mirror and from the PIC input to the loop-mirror, respectively.

With the equations above, we can get the equation of motion of the laser field:

$$\frac{dA_L}{dt} = i(\delta\omega_L - \delta\omega)A_L - \frac{\gamma_L}{2}A_L + \frac{g_L}{2}(1 + i\alpha)A_L - \frac{\kappa_L T_1 F(\delta\omega) A_L}{1 - \sqrt{T_1 R_L} F(\delta\omega)}. \quad (3.10)$$

with

$$F(\delta\omega) = T_{2c}\sqrt{R_c}e^{2i\phi_c} - \frac{\kappa_c T_{2c} e^{2i\phi_c}}{\frac{\gamma_c}{2} + i\delta\omega} - T_{2m}e^{i(\phi_{m1} + 2\phi_{m2})} \quad (3.11)$$

where $\phi_c = n_0 L_c \delta\omega / c + \phi_{c0}$ and $\phi_{m1} + 2\phi_{m2} = n_0 L_m \delta\omega / c + \phi_{m0}$. L_c is the optical path length between PIC input to the μ FP cavity, and L_m is the round-trip length after being reflected by the loop mirror from the PIC input. And n_0 is the mode index of the waveguide mode. ϕ_{c0} can be controlled by ϕ_1 and ϕ_2 , whereas $\phi_{m0} - 2\phi_{c0}$ can be controlled by ϕ_2 . It's worth noting that the first and the third term are mirror-like reflection from the μ FP cavity and the loop-mirror, whereas the middle term is the "useful" feedback for SIL. After expanding the laser field as $A_L = |A_L|e^{i\phi_L}$ and applying the stationary condition, we obtain the stationary solution of the self-

injection locked laser as:

$$0 = \delta\omega_L - \delta\omega + \alpha \frac{\gamma_L}{2} - \text{Im}[(1 + i\alpha) \frac{\kappa_L T_1 F(\delta\omega)}{1 - \sqrt{T_1 R_L} F(\delta\omega)}] \quad (3.12)$$

Typical parameters for our system are: $\alpha = 2.5$, $R_L = 1$, $R_C = 1$, $\gamma_c = 2\pi 50 \times 10^3 \text{Hz}$, $\kappa_C = 0.15\gamma_C$, $\kappa_L = 3 \times 10^5 \kappa_C$, $T_1 = 0.2$, $L_c = 0.013 \text{m}$, $n_0 = 1.47$, $L_m \approx 2L_C$. First, we examine the case where there is no reflection modification by setting $T_{2c} = 1$. This scenario mimics a laser directly coupled to a FP cavity. It is well known that such strong reflections can drive the laser into a chaotic state or introduce mode-hopping [14]. However, even if a stable external cavity mode is assumed to be achieved, this feedback significantly hinders the realization of the self-injection locking state, since it will be likely buried in the external cavity response curve (Fig. 3.8(b)).

Next, we assume perfect mirror reflection removal, where $T_{2c} = 0.5$ and $\phi_{m0} = 2\phi_{c0}$, to discuss the dynamics of SIL to the μFP cavity mode. In this case, the response curve closely resembles those observed in ring-resonator-based works [42]. The feedback phase ϕ_{c0} influences the noise-reduction factor ($\text{NRF} = \left(\frac{\partial \delta\omega_L}{\partial \delta\omega}\right)^2$) of the SIL state.

In Fig. 3.8(c), the effects of different feedback phases ϕ_{c0} are illustrated, showing an unstable SIL state (Fig. 3.8(c,i)) and stable SIL states (Fig. 3.8(c,ii-iii)) [63]. While the feedback phase ϕ_{c0} slightly impacts the noise-reduction factor, when the feedback phase remains in the stable SIL regime, its drift has minimal impact on noise performance.

Furthermore, Fig. S2(e) demonstrates how the drift of the mirror reflection removal phase, $\phi_{m0} - 2\phi_{c0}$, affects the response curve. A phase drift of $\pi/10$ impacts the response curve. However, due to the low thermo-optic coefficient of silicon nitride, we didn't observe this drift moving the laser out of the SIL state. The primary mechanism that inhibits long-term robust SIL is the drift of the DFB laser frequency,

$\delta\omega_L$. The phases used in Fig. 3(b) are $\phi_{c0} = \pi$, $\phi_{m0} - 2\phi_{c0} = 3.03478$.

We also study the dynamics of the self-injection locking by slowly modulating the DFB laser current to observe shifts in the RF beat note spectrum with a stable reference laser, as depicted in Fig. 3.9(a). This system enters and exits the SIL state at points b to b' and f to f' (Fig. 3.9(a,b)), respectively, demonstrating a locking range of 265 MHz. Comparison with a theoretical model of the self-injection locking process suggests that the onset of self-injection locking occurs earlier than expected (point b in Fig. 3.9(b) rather than the end of the curve); this is likely due to laser relaxation oscillations produced by the semiconductor gain medium [11]. The asymmetric feature for two directions of SIL is due to the fact that the feedback phase ϕ_1 is not ideal (Fig. 3.8(c)). It is also notable that away from the μ FP resonance, the relationship between output frequency and DFB laser frequency diverges from the identity line (output detuning equals to laser detuning), indicating that mirror reflections are not entirely canceled, leaving a small amount of residual broadband feedback (see SI. Fig. 3.8(e)). Nevertheless, provided that this feedback remains sufficiently weak to avoid destabilizing the laser, it appears to have little impact on the performance of the SIL laser.

3.5 Conclusion

We have demonstrated complementary approaches for the integration of high-performance vacuum-gap μ FP reference cavities with photonic integrated circuits (Fig. 3.10), to enable both electrical and optical feedback for next-generation lasers, oscillators, and sensors. Using a new circuit interface to transform the response of the μ FP resonator, we have created resonant optical feedback necessary to self-injection lock a waveguide-based laser with the reference cavity. This strategy produces an integrated laser with a fundamental linewidth of 35 mHz and an integrated $1/\pi$ linewidth of 150 Hz, cor-

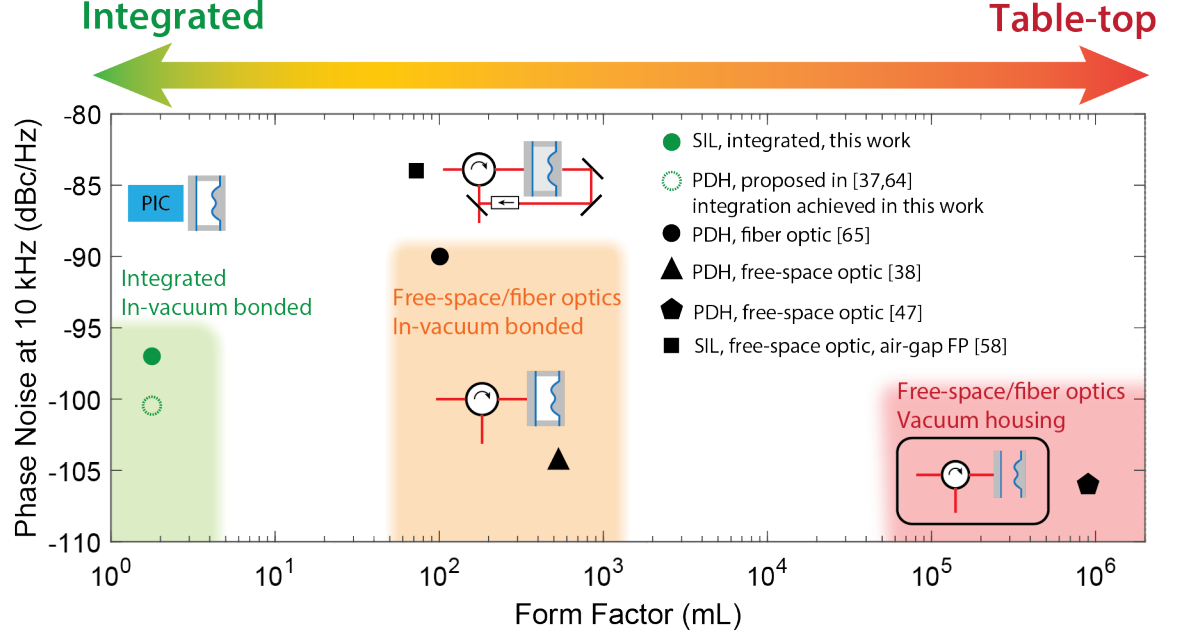


Figure 3.10: Collection of the previous results shown in [37, 38, 47, 58, 64, 65].

responding to record-level performance. Phase noise measurements represent a 19 dB (17 dB) improvement in phase noise at 1 kHz (10 kHz) offset frequencies relative to prior studies that have utilized 1.4 m spiral waveguide resonators [12]. Since the electromagnetic energy is stored almost entirely in vacuum, this vacuum-gap μ FP cavity overcomes numerous limitations imposed by dielectric-induced thermo-refractive noise. Moving from stage controlled pieces to a fully packaged and on-chip electric-optical-controlled system, we anticipate the low-offset frequency noise will be further improved and long-term locking robustness will be extended from several hours to longer durations. Such low phase noise levels make this system highly suitable for photonic microwave signal generators through use of optical frequency division, potentially eliminating the need for more complex PDH locking schemes [37].

Through a complementary approach, we have shown that the reflection response of such vacuum-gap μ FP cavities can be interrogated using a circuit interface, to enable on-chip PDH-locking and microwave photonic signal processing. Co-integration of the cavity with a silicon photonic circuit is achieved by attaching the cavity to

a multi-port vertical-emission grating coupler. Using this grating interface, the cavity is placed within an optical interferometer that cancels unwanted back-reflections, eliminating the need for on-chip circulators and isolators. This same interferometer maps the reflection response to a separate optical port, permitting efficient interrogation of the reflection response required for PDH-locking and microwave photonic signal processing. Hence, this complementary interfacing method provides access to an ultra-narrow linewidth and stable reference cavity that becomes a versatile resource for advanced microwave-photonic filtering, on-chip PDH locking, and sensing applications.

Looking beyond these specific demonstrations, the co-integration of such high-performance vacuum-gap reference cavities with photonic circuits presents numerous advantages. By storing electro-magnetic field mostly in vacuum, it mitigates the limitations associated with thermo-refractive noise and enables high performance across various photonic integrated systems. Additionally, with sealed vacuum and athermal material, these cavities significantly reduce environmental constraints, thereby broadening the operating conditions. Furthermore, such reference cavities can achieve substantially higher performance with larger mode volumes, increased cavity lengths, and improved packaging. Hence, these innovative techniques provide a crucial capabilities that are essential for the development of next-generation integrated lasers, oscillators, and sensors.

Chapter 4

On-chip Terahertz Bandwidth Nonmagnetic Isolators

4.1 Introduction

The ability to generate, manipulate, and detect light using photonic integrated circuits (PICs) of increasing complexity has opened the door to revolutionary advancements in classical [56, 66–68] and quantum computation [3, 69, 70], optical communications [71–73], as well as imaging [74–76] and sensing [77–79]. However, the realization of integrated nonreciprocal devices—such as isolators and circulators—has been challenging as they traditionally rely on magneto-optic materials [80, 81], which are incompatible with complementary metal-oxide-semiconductor (CMOS) technology. Optical isolators are needed to prevent backscattered light from creating unwanted feedback that harms the stability of lasers and amplifiers, an issue that becomes more severe due to strong scattering from waveguide sidewalls within integrated systems [15]. Backscatter also limits the number of components that can be concatenated before interference and feedback distort and destabilize the optical response to an unacceptable degree, placing a practical limit on the complexity and scale of photonic systems.

These challenges have motivated a quest for nonmagnetic forms of nonreciprocity that can be used to create broadband optical isolators within integrated photonics.

Time-modulated systems have recently been used to create nonmagnetic forms of nonreciprocity [16–18] that have yielded impressive demonstrations of optical isolation utilizing electro-optics [19,20], acousto-optics [21], and Kerr nonlinearities [22]. However, to date, constraints associated with the response of integrated resonators [22–25], the time-modulation frequency [20], or the phase-matching bandwidth of traveling wave interaction [26,27] have yielded relatively narrow ($\sim$ gigahertz) isolation bandwidths. While such narrowband isolators could be useful in specialized cases, the realization of general-purpose, wideband ($>$ terahertz) optical isolators remain critical to the integration of wideband optical amplifiers, swept lasers and broadband light sources. Moreover, by eliminating distortion and instabilities produced by unwanted feedback, wideband optical isolators permit scalable and frequency-agnostic integration of many cascaded photonic components, unlocking the true potential for integrated photonics.

Here, we combine nonreciprocal light scattering processes with wideband group delay engineering to realize a bandwidth-scalable nonmagnetic optical isolator exhibiting >1 THz isolation bandwidth. The isolator is comprised of two multi-port acousto-optic modulators that are used as beam splitters within a nonreciprocal multimode interferometer. These acousto-optic beam splitting operations are implemented using electrically-driven intermodal (inter-band) scattering process. The phase-matching condition of such scattering enables a nonreciprocal (direction-dependent) frequency shift. Within the interferometer, the direction-dependence of the frequency shift produces constructive interference for forward propagating light and cancellation for backward propagating light, yielding isolation contrast as high as 28 dB and insertion losses as low as -2.16 dB. Precise balancing of the interferometer group delay is shown to greatly enhance the isolation bandwidth, yielding isolation bandwidths

as high as 2 THz, which we demonstrate using two complementary methods. Varying the beam splitter phase using a microwave drive, we also demonstrate wideband tuning of the isolator center frequency and re-configurability of the isolator direction. Building on these results, rapidly reconfigurable isolators and circulators with much larger bandwidths ($\gg 1$ THz) should be achievable, bringing powerful new capabilities to integrated photonics.

4.2 System Concept

In this section, we present a new strategy for nonmagnetic optical nonreciprocity that enables ultra-wideband (> 1 THz) optical isolation, addressing significant challenges associated with extending isolation bandwidths. Prior works demonstrating wideband acousto-optic nonreciprocity have utilized the strict wave-vector selective nature of intermodal acousto-optic scattering to create an asymmetry between the acousto-optic scattering strengths in forward and backward directions [26, 27]. This requires distributed coupling within a long (> 1 cm) acousto-optic scattering region, yielding an isolation bandwidth that scales with the phase-matching bandwidth [26, 27]. Due to the constraints associated with phase matching and acousto-optic design, these earlier demonstrations have been limited to isolation bandwidths < 60 GHz.

To greatly extend optical isolation bandwidths, we instead harness the nonreciprocal frequency shift produced by two discrete acousto-optic modulators (AOMs) to create a wideband nonreciprocal interferometer. In contrast to these prior works [26, 27], our current scheme uses a very short (< 300 μm) acousto-optic interaction length with relaxed wave-vector selectivity. This enables a direction-dependent frequency shift across a large wavelength range, permitting wideband isolator performance. Each AOM is used to split an incoming wave into two wave components having distinct frequencies and spatial modes, which enter or exit the AOM through different ports

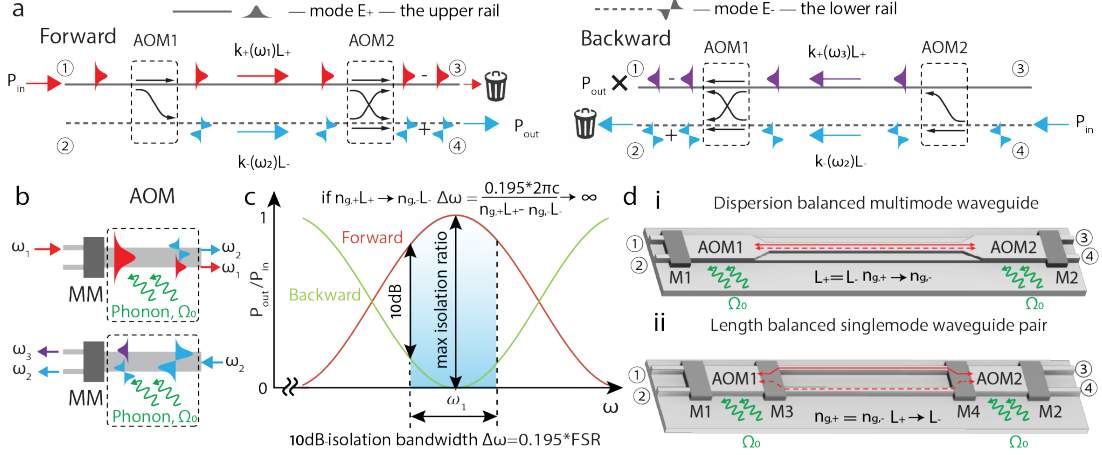


Figure 4.1: **a** Dual-rail representation of the bandwidth-scalable isolator for forward (left) and backward (right) operation. AOM1 and AOM2 denote acousto-optic modulators 1 and 2, respectively, acting as nonreciprocal beam splitters. **b** The nonreciprocal acousto-optic modulator in our system partially scatters the forward-propagating symmetric mode at frequency ω_1 into an antisymmetric mode at frequency $\omega_2 = \omega_1 + \Omega_0$, while partially scattering the backward-traveling antisymmetric mode at frequency ω_2 into a symmetric mode at frequency $\omega_3 = \omega_2 + \Omega_0$. Here Ω_0 represents phonon angular frequency. A mode-multiplexer (MM) is used to selectively excite and extract different spatial mode in this multimode waveguide. For detailed discussion on AOM nonreciprocity, please refer to SI. Sec I. **c** Theoretical response of the nonreciprocal multimode interferometer. Similar to a conventional Mach-Zehnder interferometer, the response exhibits a cosine-shaped dependence on optical frequency. The nonreciprocal beam splitter—implemented here by the AOM—shifts the forward and backward responses, thus functioning as an isolator. By balancing the two arms of the interferometer, characterized by $n_{g,+}L$ and $n_{g,-}L$, the isolation bandwidth is increased. **d** Two complementary approaches for bandwidth-scalable optical isolators. Optical paths in the delay region for symmetric mode and anti-symmetric mode are denoted by solid and dashed line respectively, same as panel **a**. **i** Balancing the interferometer by transiting the regular multi-mode waveguide to a dispersion balanced ($n_{g,+} \cong n_{g,-}$) multimode waveguide. **ii** Balancing the interferometer by mapping two modes into two identical waveguides ($L_+ \cong L_-$) through mode multiplexers (M3, M4).

of a spatial mode multiplexer (Fig. 4.1b). Using two AOMs as frequency-shifting beamsplitters, we create a Mach-Zehnder interferometer that uses wave interference to produce nonreciprocal light transmission (see Fig. 4.1a). By precisely balancing the phase delay of the arms of this interferometer, we show that the isolation bandwidths of this system are readily extended to > 1 THz. In the following, We present two

complementary implementations of this wideband isolator. The first implementation (Fig. 4.1d(i)) uses a dispersion-balanced multi-mode waveguide to achieve phase- and group-delay balancing, while the second (Fig. 4.1d(ii)) uses distinct optical waveguides to achieve balancing.

Both experimental implementations, seen in Fig. 4.1d(i) and Fig. 4.1d(ii), are realized using silicon photonic circuits. Each AOM uses electrically transduced traveling-wave phonons through an interdigital transducer (IDT) to produce phase-matched intermodal scattering within a multimode silicon waveguide segment [27]. This acousto-optic interaction transfers energy between a symmetric and antisymmetric TE-like waveguide modes, imparting a direction-dependent frequency shift to the scattered waves (Fig. 4.1b). As illustrated in Fig. 4.1d and Fig. 4.1b, spatial mode multiplexers are used to selectively address the symmetric ($\mathbf{E}_+$) and antisymmetric ($\mathbf{E}_-$) waveguide modes at the input and output of the acousto-optic waveguide segments. Through both implementations (see Fig. 4.1d) two acousto-optic modulators are cascaded in series to close the paths of the interferometer. In the first implementation (Fig. 4.1d(i), device A and B)), phase- and group-delay balancing between the two optical paths is implemented using a dual-ridge dispersion-engineered multimode waveguide. Through the second implementation (Fig. 4.1d(ii), device C), balancing is achieved using distinct single mode waveguides that are interfaced with two additional mode multiplexers (M2, M3).

4.2.1 Electrically driven acousto-optic scattering

Before delving into the detailed operating principle of our ultra-wideband isolator, we first review the fundamentals of intermodal acousto-optic scattering. We then illustrate how integrated AOMs effectively function as nonreciprocal beam splitters.

We assume that the angled AOM launches traveling-wave phonons with a frequency of $\Omega/2\pi$ and a wavevector of $q(\Omega)$ along the waveguide. In order to achieve

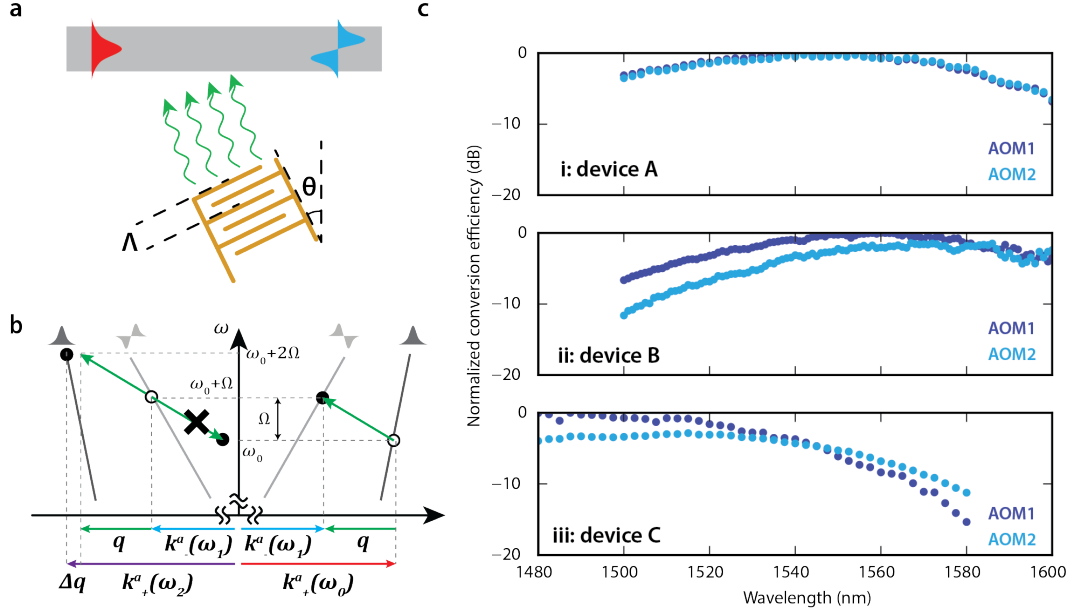


Figure 4.2: **a**, Schematic of single section electro-optomechanical scattering. **b**, Dispersion diagram of optomechanical scattering. **c**, Single section measurement of three different devices with a RF drive at 2.36 GHz. **i**, **ii** refer to dispersion-balanced multimode waveguide device A and B **iii** refers to length balanced singlemode waveguide device.

effective forward operation of the isolator via acoustic optical scattering, it is crucial that the condition for phase matching is met. This can be mathematically expressed as:

$$k_+(\omega_0) + q(\Omega) = k_-(\omega_0 + \Omega) \quad (4.1)$$

In this equation, $k_+(\omega_0)$ and $\omega_0/2\pi$ represent the wavevector and frequency of the incident light in the symmetric mode, respectively, while $k_-(\omega_0 + \Omega)$ is the wavevector of the anti-Stokes scattered light in the anti-symmetric mode, as they are shown in Fig. 4.2(b).

If we examine a scenario where the incident optical frequency is slightly off the ideal phase matching condition, a phase mismatch will inevitably occur. This phase mismatch can be quantitatively evaluated using the following equation:

$$\Delta q_{\text{pm}}(\omega) = k_+(\omega) + q(\Omega) - k_-(\omega + \Omega) = \frac{n_{\text{g},+} - n_{\text{g},-}}{c}(\omega - \omega_0) \quad (4.2)$$

In this equation, $n_{\text{g},1(2)}$ signifies the group index of the symmetric (anti-symmetric) modes within the active waveguide.

Following the establishment of the phase mismatch, the frequency response of modulation will bear a proportional relationship to $|\int_0^{L_a} \exp i\Delta q_{\text{pm}} z dz|^2 = L_a^2 \text{sinc}^2(\Delta q_{\text{pm}} L_a/2)$, with L_a being the interaction length within the active region. Given our interaction length L_a is relatively short, $\Delta q_{\text{pm}} L_a$ will be small over a broadband of optical wavelength. And even in backward direction, we have a phase mismatch term $\Delta q_{\text{pm}} = 2\frac{n_{\text{g},+}}{c}\Omega$, as is shown in Fig. 4.2(b), the anti-Stokes scattering can still happen since $\Delta q_{\text{pm}} L_a$ is small. However the Stokes process won't happen in backward operation because of phase mismatch being too large. Such phase matching condition provides us a non-reciprocal acoustic optical modulation which is essential to make an isolator. In this way, such a short AOM can produce a direction dependent frequency shift, making it a non-reciprocal beam splitter.

It is important to note that the acoustic wavevector is determined by the IDT tooth period (Λ) and angle (θ), as given by:

$$q(\Omega) = \frac{2\pi}{\Lambda} \sin \theta \quad (4.3)$$

Therefore, based on equation above, the center frequency of conversion will be influenced by Λ, θ parameters of the IDT. Here our IDT tooth period is fixed at $3.5\mu\text{m}$, and IDT angle is different for three devices we showed in the main text. The single section conversion efficiency characterization is shown in Fig. 4.2(c). Device A, B and C have the IDT angle of 8.59° , 8.77° and 7.94° , showing different center wavelength and broadband response.

4.2.2 Cavity enhanced acoustic-optical scattering

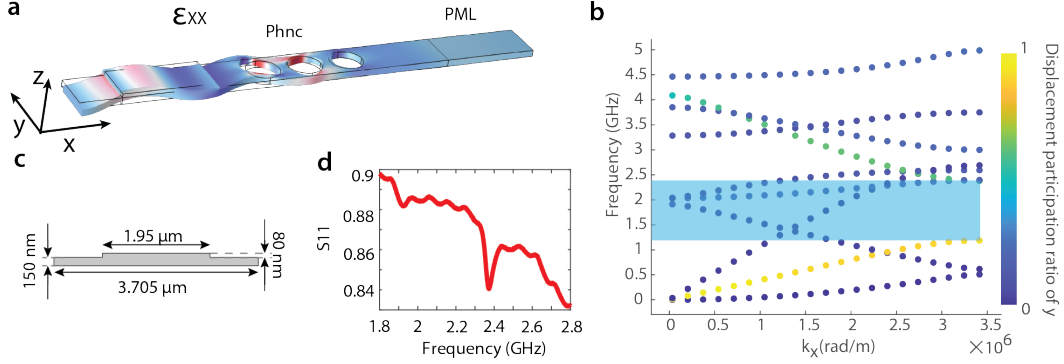


Figure 4.3: **a**, Acoustic cavity simulation for the mode favorable for inter-mode scattering. **b**, Band diagram of phononic crystal. **c**, Dimensions for the multimode waveguide of AOM active region. **d**, S_{11} measurement of the IDT.

To improve acousto-optic scattering efficiency, we employ an acoustic cavity to enhance the acoustic field. The acoustic cavity design, depicted in Fig.4.3(a), is a Fabry-Pérot cavity consisting of a free-boundary mirror on one side and a phononic crystal mirror on the other. The phononic crystal, formed by etching 710 nm diameter holes in silicon with a 920 nm period, generates a stop band ranging from 1.25 GHz to 2.38 GHz, as shown in Fig.4.3(b). Simulations indicate that the acoustic cavity achieves a quality factor of 512 at 2 GHz, primarily limited by leakage through the 3-row phononic crystal mirror, which matches well with a critically coupled cavity of total $Q=230$.

In practical implementations, however, measurements revealed a resonance frequency of 2.36 GHz, likely shifted due to stress introduced by the AlN layer. The S_{11} response of the IDT, optimized for our AOM, is illustrated in Fig. 4.3(c). The response demonstrates a linewidth of approximately 50 MHz centered at 2.37 GHz.

The microwave power efficiency for electro-optomechanical scattering can be further optimized by several simple design optimization. First we can suspend the IDT such that most of the acoustic energy is confined in Si layer rather than in the bottom silica cladding of the device layer [26]. Secondly, we can tune phononic crystal di-

mensions to better match the stopgap to IDT response. Thirdly, we can improve the impedance matching of IDT to $50\ \Omega$ by increasing the IDT area. It will also improve the acoustic optical coupling for longer interaction length.

4.2.3 Extending isolation bandwidth

Next, we introduce the operating principle shared by both isolator designs. For clarity, we use a simplified dual-rail representation (Fig. 4.1a-b) to artificially separate the optical path for each mode within this system. In the forward direction (Fig. 4.1a), light of frequency ω_1 enters Port 1 to excite the symmetric mode at frequency ω_1 , denoted as $\mathbf{E}_+(\omega_1)$. As this incident wave traverses AOM 1, a portion of light is blue-shifted to $\omega_2 = \omega_1 + \Omega_0$ and transferred to the antisymmetric mode $\mathbf{E}_-(\omega_2)$ (see Fig. 4.1b). Here, Ω_0 is the angular frequency of the phonon modes. With propagation lengths of L_+ and L_- for the symmetric and antisymmetric modes, these forward-propagating wave components acquire a phase difference $\phi_f = k_+(\omega_1)L_+ - k_-(\omega_2)L_-$ before entering AOM 2. As these wave components traverse AOM 2, part of the light in mode $\mathbf{E}_-(\omega_2)$ is scattered to $\mathbf{E}_+(\omega_1)$, and part of the light in mode $\mathbf{E}_+(\omega_1)$ is scattered to $\mathbf{E}_-(\omega_2)$, resulting in coherent interference of the transmitted and scattered waves in each spatial mode. With appropriate choice of passive waveguide length (L_+ , L_-) and phase difference (ϕ_f), we produce destructive interference for waves scattered by AOM 2 into mode $\mathbf{E}_+(\omega_1)$ and constructive interference in mode $\mathbf{E}_-(\omega_2)$, resulting in light transmission from Port 4 of the system.

Due to the nonreciprocal nature of inter-modal acousto-optic scattering, waves entering Port 4 do not retrace the same path in the backward direction. A reflected wave of frequency ω_2 that enters AOM 2 through mode $\mathbf{E}_-(\omega_2)$, is blue-shifted to frequency $\omega_3 = \omega_2 + \Omega_0$ and scattered to mode $\mathbf{E}_+(\omega_3)$. In contrast to the forward-propagating case, the light in mode $\mathbf{E}_+(\omega_3)$ is higher in frequency than that in mode $\mathbf{E}_-(\omega_2)$. Hence, wave components $\mathbf{E}_+(\omega_3)$ and $\mathbf{E}_-(\omega_2)$ accumulate a distinct relative

phase ($\phi_b = k_+(\omega_3)L_+ - k_-(\omega_2)L_-$) as they propagate through the passive waveguide section in the backward direction; consequently, they do not interfere the same way when recombined, giving rise to nonreciprocal light transmission. Optical isolation is produced when paths of this multimode system are designed to simultaneously produce constructive interference in the forward direction ($\phi_f = 0$) and destructive interference in the backward direction ($\phi_b = \pm\pi$). To satisfy these conditions, the nonreciprocal phase difference, $\Delta\phi$, must meet the requirement $\Delta\phi = \phi_f - \phi_b = \pm\pi$. Since $\Delta\phi$ can be expressed as $\Delta\phi = k_+(\omega_1)L_+ - k_+(\omega_3)L_+ = -2v_{g,+}\Omega_0L_+$, we see that this condition for nonreciprocal light transmission is directly controlled by the delay length L_+ , the group velocity for the symmetric mode $v_{g,+}$, and the frequency shift, Ω_0 , produced by the AOM.

The isolation bandwidth is determined by the difference in effective group delay between the interferometer paths, analogous to the free spectral range (FSR) of a Mach-Zehnder interferometer. Fig. 4.1c shows the total power transmissivity between Port 1 and Port 4 of the nonreciprocal multimode interferometer as a function of the incident optical frequency. By carefully tuning the nonreciprocal phase, $\Delta\phi$, the sinusoidal forward (red) and backward (green) optical response become out of phase, resulting in wideband optical isolation centered around ω_1 . We use the 10 dB isolation bandwidth, the bandwidth over which at least 10 dB isolation contrast is achieved, to specify the isolator performance. The 10 dB isolation bandwidth, $\Delta\omega$, is inversely proportional to the group delay difference and is given by $\Delta\omega = 0.195 \times \text{FSR} = 0.195 \times 2\pi c / (n_{g,+}L_+ - n_{g,-}L_-)$, where c is the speed of light and $n_{g,+}$ and $n_{g,-}$ are the group indices of the symmetric and antisymmetric modes, respectively. (See supplement Sec. I for details.) Therefore, by balancing the group delay of the interferometer, we can theoretically extend the isolation bandwidth without limit.

In what follows, we describe two distinct methods for wideband group delay balancing. The first method utilizes a multimode two-ridge waveguide, where two optical

modes traverse the same distance of length $L = L_+ = L_-$ (Fig. 4.1d(i)). The geometry of the separate waveguide ridges are individually tuned to control the group indices of the two optical modes, such that $n_{g,+} \cong n_{g,-}$. The second method, alternatively, features the design of a waveguide pair, demultiplexing the two optical modes into separate identical single-mode waveguides (Fig. 4.1d(ii)), with group index $n_{g,0}$ such that $n_{g,+} = n_{g,-} = n_{g,0}$. In this case, we balance group delay of two paths by independently adjusting the waveguide lengths L_+ and L_- to realize wideband isolation.

Transfer matrix analysis

In the next, we present a detailed analysis of our nonreciprocal multimode interferometer, presenting a way to explain its operation principles (Fig. 4.1(a)). For forward operation, we describe the system's response as follows

$$\begin{bmatrix} a_{o,+, \omega_1} \\ a_{o,-, \omega_1} \end{bmatrix} = t_a^f(\phi_{\text{RF2}}) \cdot t_{\text{delay}} \cdot t_a^f(\phi_{\text{RF1}}) \begin{bmatrix} a_{i,+, \omega_1} \\ a_{i,-, \omega_2} \end{bmatrix}. \quad (4.4)$$

Here, the transfer matrices $t_a^f(\phi_{\text{RF1(2)}})$ represent the first and second AOM section with input RF phase $\phi_{\text{RF1(2)}}$, respectively, while t_{delay} accounts for the delay introduced between these AOMs.

Neglecting waveguide losses, the transfer matrix for the AOM section can be expressed as [27]

$$t_a^f(\phi_{\text{RF}}) = \begin{bmatrix} e^{[i(k_+ - \frac{\Delta q_{\text{pm}}}{2})L_a]} \left\{ \cos[\beta L_a] + \frac{i\Delta q_{\text{pm}}}{2\beta} \sin[\beta L_a] \right\} & -e^{[i(k_+ - \frac{\Delta q_{\text{pm}}}{2})L_a]} \frac{ig^* \bar{b}^*}{\sqrt{v_1 v_2}} \sin[\beta L_a] \\ -e^{[i(k_- + \frac{\Delta q_{\text{pm}}}{2})L_a]} \frac{ig\bar{b}}{\sqrt{v_1 v_2}} \sin[\beta L_a] & e^{[i(k_- + \frac{\Delta q_{\text{pm}}}{2})L_a]} \left\{ \cos[\beta L_a] - \frac{i\Delta q_{\text{pm}}}{2\beta} \sin[\beta L_a] \right\} \end{bmatrix}. \quad (4.5)$$

In this equation, $\beta \equiv \sqrt{\left(\frac{\Delta q_{\text{pm}}(\omega)}{2}\right)^2 + \frac{|\bar{b}g|^2}{v_+v_-}}$, with $v_{+(-)}$ denoting the group velocities of the symmetric and anti-symmetric modes in AOM region, and $\bar{b}$ and g signifying the acoustic field and acoustic-optical coupling strength, respectively.

To simplify the transfer matrix, we introduce the definitions: $r(\omega)e^{i\phi_r(\omega)} \equiv \cos[\beta L_a] + \frac{i\Delta q_{\text{pm}}}{2\beta} \sin[\beta L_a]$, $\mu(\omega)e^{i\phi_{\text{RF}}} \equiv \frac{g\bar{b}}{\sqrt{v_1v_2}} \sin[\beta L_a]$, where $r(\omega), \mu(\omega)$ are real and $r(\omega)^2 + \mu(\omega)^2 = 1$. $\phi_r(\omega)$ can be understood as the phase originated from acoustic optical phase mismatch, and ϕ_{RF} is the acoustic phase that can be controlled by the driving RF phase.

$$t_a^f(\phi_{\text{RF}}) = \begin{bmatrix} r(\omega)e^{i(k_+ - \frac{\Delta q_{\text{pm}}}{2})L_a}e^{i\phi_r(\omega)} & -i\mu(\omega)e^{i(k_+ - \frac{\Delta q_{\text{pm}}}{2})L_a}e^{-i\phi_{\text{RF}}} \\ -i\mu(\omega)e^{i(k_- + \frac{\Delta q_{\text{pm}}}{2})L_a}e^{i\phi_{\text{RF}}} & r(\omega)e^{i(k_- + \frac{\Delta q_{\text{pm}}}{2})L_a}e^{-i\phi_r(\omega)} \end{bmatrix}. \quad (4.6)$$

The transfer matrix for the delay region is expressed as

$$t_{\text{delay}} = \begin{bmatrix} e^{ik_{d,+}(\omega_1)L_{d,+}} & 0 \\ 0 & e^{ik_{d,-}(\omega_2)L_{d,-}} \end{bmatrix}, \quad (4.7)$$

where $k_{d,+(-)}, L_{d,+(-)}$ are the wavevector and delay length in delay region for symmetric mode and anti-symmetric mode.

For forward operation, given only input from $a_{i,+, \omega_1} = a_0$, we can calculate the output as

$$\begin{bmatrix} a_{o,+, \omega_1} \\ a_{o,-, \omega_1} \end{bmatrix} = a_0 \begin{bmatrix} r(\omega_1)^2 e^{i[(2k_+ - \Delta q_{\text{pm}})L_a + 2\phi_r + k_{d,+}L_{d,+}]} - \mu(\omega_1)^2 e^{i[(k_+ + k_-)L_a + k_{d,-}L_{d,-} + \phi_{\text{RF}1} - \phi_{\text{RF}2}]} \\ -ir(\omega_1)\mu(\omega_1)(e^{i[(k_- + k_+)L_a + \phi_{\text{RF}2} + \phi_r + k_{d,+}L_{d,+}]} + e^{i[(2k_- + \Delta q_{\text{pm}})L_a - \phi_r + \phi_{\text{RF}1} + k_{d,-}L_{d,-}]}) \end{bmatrix}. \quad (4.8)$$

Similarly, for backward operation, we have

$$\begin{bmatrix} a_{o,+,\omega_3} \\ a_{o,-,\omega_1} \end{bmatrix} = t_a^b(\phi_{\text{RF1}}) \cdot t_{\text{delay}} \cdot t_a^b(\phi_{\text{RF2}}) \begin{bmatrix} a_{i,+,\omega_3} \\ a_{i,-,\omega_2} \end{bmatrix}. \quad (4.9)$$

The only difference between t_a^f and t_a^b is that t_a^b has an extra phase mismatch term $q_{\text{pm}}^b = q_{\text{pm}}^f + 2\frac{n_{\text{g},+}}{c}\Omega$. Since our interaction region is short, $2\frac{n_{\text{g},+}}{c}\Omega L_a \ll 1$, this term can be disregarded. Then we have $t_a^f \approx t_a^b$.

Given only input from $a_{i,-,\omega_2} = a_0$, the output for backward operation can be calculated as

$$\begin{bmatrix} a_{o,+,\omega_3} \\ a_{o,-,\omega_1} \end{bmatrix} = a_0 \begin{bmatrix} -ir(\omega_3)\mu(\omega_3)(e^{i[(k_-+k_+)L_a+\phi_{\text{RF1}}+\phi_r+k_{d,-}L_{d,-}]} + e^{i[(2k_++\Delta q_{\text{pm}})L_a-\phi_r+\phi_{\text{RF2}}+k_{d,+}L_{d,+}]}) \\ r(\omega_3)^2 e^{i[(2k_--\Delta q_{\text{pm}})L_a+2\phi_r+k_{d,-}L_{d,-}]} - \mu(\omega_3)^2 e^{i[(k_++k_-)L_a+k_{d,+}L_{d,+}+\phi_{\text{RF1}}-\phi_{\text{RF2}}]} \end{bmatrix}. \quad (4.10)$$

For simplicity, we can assume small AOM conversion efficiency, $\mu(\omega) \ll 1$, to study the frequency response of the system. With this assumption, we can write simplified output for forward operation as

$$\begin{bmatrix} a_{o,+,\omega_1} \\ a_{o,-,\omega_1} \end{bmatrix} = a_0 \begin{bmatrix} e^{i(2k_++k_{d,+}L_{d,+})} \\ -ir(\omega_1)\mu(\omega_1)(e^{i[(k_-+k_++\Delta q_{\text{pm}}/2)L_a+\phi_{\text{RF2}}+k_{d,-}L_{d,-}]} + e^{i[(2k_++\Delta q_{\text{pm}}/2)L_a+\phi_{\text{RF1}}+k_{d,+}L_{d,+}]}) \end{bmatrix}. \quad (4.11)$$

Here we have used the fact that when $\mu(\omega) \ll 1$, $2\phi_r = \Delta q_{\text{pm}}L_a$.

Similarly, the simplified output for backward operation is

$$\begin{bmatrix} a_{o,+,\omega_3} \\ a_{o,-,\omega_1} \end{bmatrix} = a_0 \begin{bmatrix} -ir(\omega_3)\mu(\omega_3)(e^{i[(k_-+k_++\Delta q_{pm}/2)L_a+\phi_{RF1}+k_{d,+}L_{d,+}]} + e^{i[(2k_-+\Delta q_{pm}/2)L_a+\phi_{RF2}+k_{d,-}L_{d,-}]}) \\ e^{i(2k_-+k_{d,-}L_{d,-})} \end{bmatrix} \quad (4.12)$$

In order to design an isolator with minimal insertion loss, we aim to maximize $a_{o,-,\omega_1}$ in forward operation and simultaneously minimize $a_{o,+,\omega_3}$ in backward operation. To achieve this, we derive the phase relations that satisfy these requirements

$$(k_+(\omega_1) - k_-(\omega_2))L_a + k_{d,+}(\omega_1)L_{d,+} - k_{d,-}(\omega_2)L_{d,-} + \phi_{RF1} - \phi_{RF2} = 2n\pi, \quad (4.13)$$

$$(k_+(\omega_3) - k_-(\omega_2))L_a + k_{d,+}(\omega_3)L_{d,+} - k_{d,-}(\omega_2)L_{d,-} + \phi_{RF1} - \phi_{RF2} = (2m-1)\pi. \quad (4.14)$$

Here, n and m are integers. For our analysis, we consider the designed central optical frequency $\omega_0/2\pi$, with $\omega_1 = \omega_0$, $\omega_3 = \omega_2 + \Omega = \omega_1 + 2\Omega$, and $\phi_{RF1} = \phi_{RF2}$. Using (S.13) – (S.14), we arrive at the minimum phase delay requirement (where $n = m$) as

$$\frac{2\Omega n_{g,+}L_a}{c} + \frac{2\Omega n_{g,+}^d L_{d,+}}{c} = \pi. \quad (4.15)$$

This equation provides the design guidelines of the delay line length for two dispersion engineering methods, which we will discuss in more details in the following sections.

Dispersion balanced multimode waveguide

Here we adapt the theory derived in above to the dispersion engineered waveguide method. Following the equation S.13-14, the forward and backward can be viewed as the interference of two scattering terms. Then the anti-Stokes scattering output power in forward and backward operation can be written as

$$P_{as}^f \propto |a_{as,0}|^2 \cos^2 \frac{\phi_f(\omega)}{2}; \quad (4.16)$$

$$P_{\text{as}}^{\text{b}} \propto |a_{\text{as},0}|^2 \cos^2 \frac{\phi_{\text{b}(\omega)}}{2}, \quad (4.17)$$

where $|a_{\text{as},0}|^2 = |g\bar{b}|^2 L_{\text{a}}^2 \text{sinc}^2(\Delta q_{\text{pm}}(\omega) L_{\text{a}}/2)$ and $\phi_{\text{f}(\text{b})(\omega)}$ is phase difference between two terms for forward (backward) direction. In dispersion engineered waveguide method, we have $L_{\text{d},+} = L_{\text{d},-} = L_{\text{d}}$, and using Taylor expansion, we have

$$\begin{aligned} \phi_{\text{f}}(\omega) &= (k_+(\omega) - k_-(\omega + \Omega))(L_{\text{a}} + L_{\text{t}}) + (k_+^{\text{d}}(\omega) - k_-^{\text{d}}(\omega + \Omega))L_{\text{d}} + \phi_{\text{RF1}} - \phi_{\text{RF2}} \\ &= q(L_{\text{a}} + L_{\text{t}}) + \phi_{\text{RF1}} - \phi_{\text{RF2}} + \frac{n_{\text{g},+} - n_{\text{g},-}}{c}(\omega - \omega_0)(L_{\text{a}} + L_{\text{t}}) \\ &\quad + \frac{n_{\text{g},+}^{\text{d}} - n_{\text{g},-}^{\text{d}}}{c}(\omega - \omega_0)L_{\text{d}} + \frac{\beta_{\text{d},+} - \beta_{\text{d},-}}{2}(\omega - \omega_0)^2 L_{\text{d}} \end{aligned} \quad (4.18)$$

$$\begin{aligned} \phi_{\text{b}}(\omega) &= (k_+(\omega + \Omega) - k_-(\omega))(L_{\text{a}} + L_{\text{t}}) + (k_+^{\text{d}}(\omega + \Omega) - k_-^{\text{d}}(\omega))L_{\text{d}} + \phi_{\text{RF1}} - \phi_{\text{RF2}} \\ &= \phi_{\text{f}}(\omega) + \left(\frac{n_{\text{g},+} + n_{\text{g},-}}{c}(L_{\text{a}} + L_{\text{t}}) + \frac{n_{\text{g},+}^{\text{d}} + n_{\text{g},-}^{\text{d}}}{c}L_{\text{d}} \right) \Omega \end{aligned} \quad (4.19)$$

where $n_{\text{g},+(-)}^{\text{d}}, \beta_{\text{d},+(-)}$ are the group index and group velocity dispersion for symmetric (anti-symmetric) mode. It's worth to mention that such dispersion engineered waveguide is able to balance $n_{\text{g},+}$ and $n_{\text{g},-}$, but for $\beta_{\text{d},+(-)}$, they are not guaranteed to be balanced. In such case we need to take this higher order term into consideration. Here for device B we design second order term, $\frac{\beta_{\text{d},+} - \beta_{\text{d},-}}{2}(\omega - \omega_0)^2 L_{\text{d}}$, to be cancelling with first order terms at the wavelength we're interested, such that we have flat response and ultra wide-band isolation (Fig. 4.4). And we also introduce L_{t} to account for the transition between the AOM and delay region.

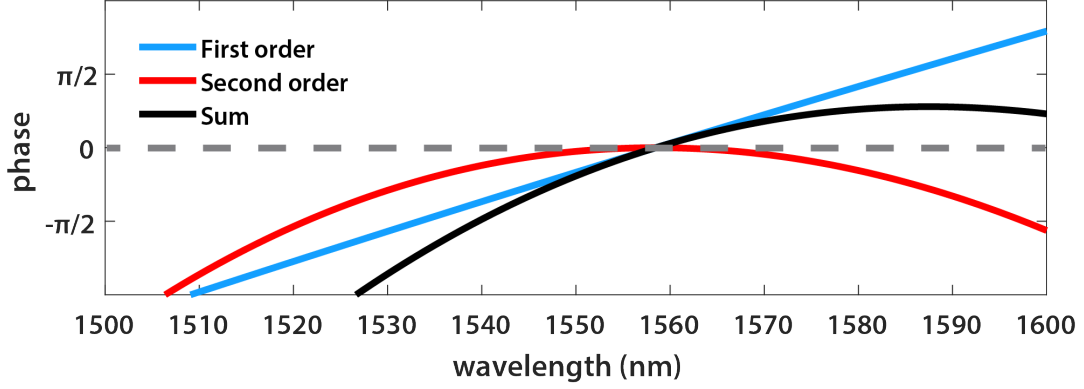


Figure 4.4: Illustration of the contribution from different group delay terms using Device B fitted data.

Length balanced single mode waveguide pair

The primary difference between the two-waveguide and two-ridge approaches lies in the phase delay region. In the two-waveguide method, the phase delay can be expressed as follows:

For the forward direction, we have

$$\begin{aligned}
 \phi_f(\omega) = & (k_+^a(\omega) - k_-^a(\omega + \Omega))L_a + (k^d(\omega)L_{d,+} - k^d(\omega + \Omega)L_{d,-}) + \phi_{\text{RF1}} - \phi_{\text{RF2}} \\
 = & q(L_a + L_t) + \phi_{\text{RF1}} - \phi_{\text{RF2}} + \frac{n_{g,+} - n_{g,-}}{c}(\omega - \omega_0)L_a + \frac{n_g^d L_{d,+} - n_g^d L_{d,-}}{c}(\omega - \omega_0) \\
 & + \frac{\beta_d}{2}(\omega - \omega_0)^2(L_{d,+} - L_{d,-}).
 \end{aligned} \tag{4.20}$$

And for the backward direction, we have

$$\begin{aligned}
 \phi_b(\omega) = & (k_+^a(\omega + \Omega) - k_-^a(\omega))L_a + (k^d(\omega + \Omega)L_{d,+} - k^d(\omega)L_{d,-}) + \phi_{\text{RF1}} - \phi_{\text{RF2}} \\
 = & \phi_f(\omega) + \left(\frac{n_{g,+} + n_{g,-}}{c}L_a + \frac{n_g^d L_{d,+} + n_g^d L_{d,-}}{c} \right) \Omega.
 \end{aligned} \tag{4.21}$$

Here, $k_d(\omega)$, n_g^d and β_d denote the wavevector, the group index and the group veloc-

ity dispersion of the single-mode waveguide in the phase delay region, respectively. In contrast to dispersion engineering, tuning waveguide length can simultaneously balance group delay difference and higher order delay originated from group velocity dispersion. In this case we can neglect this second order effect of dispersion properties, making our design much simpler.

In practical scenarios, we cannot always assume that $L_{d,+}$ is significantly greater than L_a . Ideally, our goal is to satisfy the following equation:

$$(n_{g,+} - n_{g,-}) L_a + n_g^d (L_{d,+} - L_{d,-}) = 0, \quad (4.22)$$

which would allow for the expansion of the isolation bandwidth to infinity. Achieving this condition can be straightforward in the two-waveguide method, by adjusting the length difference between the upper and lower single-mode waveguides such that $L_{d,+} - L_{d,-} = -\frac{n_{g,+} - n_{g,-}}{n_g^d} L_a$.

4.3 Experimental Results

4.3.1 Terahertz bandwidth isolation based on dispersion balanced multimode waveguide

We begin by developing the isolator that utilizes a dispersion-balanced multimode waveguide, as shown by the device microscopy presented in the inset of Fig. 4.5a. Both dispersion engineering wideband isolators were fabricated in an AlN-on-silicon platform using a CMOS compatible process. The AOMs use an angled interdigitated transducer (IDT) patterned on a piezoelectric AlN layer atop the silicon layer to excite elastic waves with the appropriate wavevector and frequency, which resonantly tunnel into the silicon multimode waveguide through a partially-transmitting phononic crystal (zoom-in of Fig. 4.5a). A membrane suspended optomechanical waveguide is

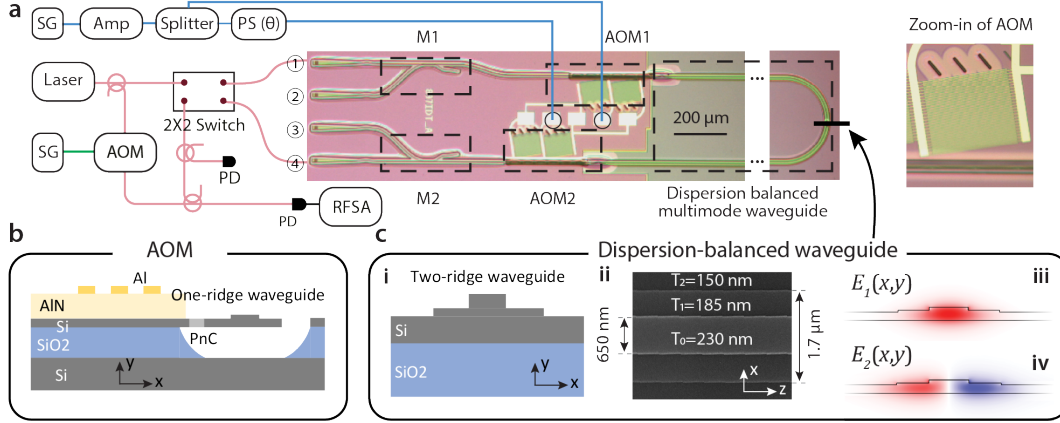


Figure 4.5: **a**, The isolator microscopy and experimental setup for characterizing the forward and backward operation of our isolator. SG: signal generator; Amp: RF amplifier; Splitter: 50/50 power splitter; PS: RF phase shifter; AOM: acousto-optic modulator; PD: photodiode; RFSa: RF spectral analyzer; M1/2: mode-multiplexer. Pink lines represent optical connections, blue lines represent high frequency ($\sim 2\text{GHz}$) electrical connections and green lines represent low frequency (55MHz) electrical connections. **b**, The on-chip acousto-optic modulator (AOM) comprises a piezoelectric interdigital transducer (IDT) fabricated on AlN and an optomechanical waveguide crafted from a suspended single silicon ridge. **c**, The dispersion-balanced waveguide is an unsuspended two-ridge waveguide. A SEM image is shown in **ii** and the designed dimensions of the two-ridge waveguide are labeled. Such waveguide provides similar group index of the symmetric (**iii**) and antisymmetric (**iv**) modes.

designed to produce simultaneous guidance of both optical and acoustic waves, facilitating efficient acousto-optic scattering (Fig. 4.5b) [27]. Optical waves are confined through total internal reflection within the central ridge while acoustic waves are confined by a phononic crystal and an air slot on either side of the ridge. To selectively excite and read out a specific spatial mode within the multimode waveguide of the AOM, we employ two mode-multiplexers (M1 and M2, as shown in Fig. 4.5a).

In this system, the group delay experienced by the symmetric ($\mathbf{E}_+$) and antisymmetric ($\mathbf{E}_-$) waveguide modes are precisely balanced by optimizing the dimensions of the two-ridge multimode waveguide seen in Fig. 4.5c.i and ii. The narrow and wide silicon ridges control the dispersion of the symmetric (Fig. 4.5c.iii) and antisymmetric (Fig. 4.5c.iv) modes, respectively; the chosen waveguide dimensions (see Fig. 4.5c.ii) equalize of the group delay of the $\mathbf{E}_+$ and $\mathbf{E}_-$ modes, which define the two paths of the

interferometer (see in Supplementary Information Sec.II (B-C) for further details). To meet the necessary conditions for constructive (destructive) interference in the forward (backward) direction, $-2v_{g,+}\Omega_0L = -\pi$, a dispersion-engineered waveguide of length $L = 7.64$ mm is used to link AOM 1 and AOM 2 within the interferometer. The performance of the AOMs within the isolator are analyzed using the experimental setup of Fig. 4.5a. The scattering efficiency and spectral response of each AOM is analyzed by mixing the transmitted optical waves with an electrically synthesized optical local oscillator (LO) while measuring the detected heterodyne spectrum. The measured RF-frequency-dependent conversion efficiencies of AOM 1 and AOM 2 are shown in Fig. 4.6a.i at a constant microwave drive power of 32.39 dBm. The conversion efficiency is defined as the fraction of incident laser power entering port 1 that is blue-shifted before exiting the system through port 4. Measurements reveal a peak acousto-optic scattering efficiency of -4.6 dB at a ~ 2.36 GHz drive frequency; this peak in conversion efficiency identifies the frequency at which phonons emitted by the IDT excite the guided acoustic wave resonance, calculated a Q-factor of ~ 230 through Lorentz fitting.

When both AOM 1 and AOM 2 are electrically driven, transmission measurements reveal constructive (destructive) interference of the acousto-optically scattered wave components in the forward (backward) direction, producing high-contrast optical isolation. To analyze the response of the system, we again perform frequency-resolved transmission measurements (Fig. 4.6a.ii), which show the portion of incident laser power that is blue-shifted by the AOMs as it traverses the system. To examine the forward operation of the isolator, we inject 1547.5 nm light into Port 1 and measure the frequency-shifted optical transmission from Port 4 while sweeping the drive frequency of AOM 1 and AOM 2. At an RF drive of 2.36 GHz, frequency-resolved conversion efficiency measurements in the forward direction (blue) show a 6 dB increase compared to a single AOM, demonstrating the coherent summation of

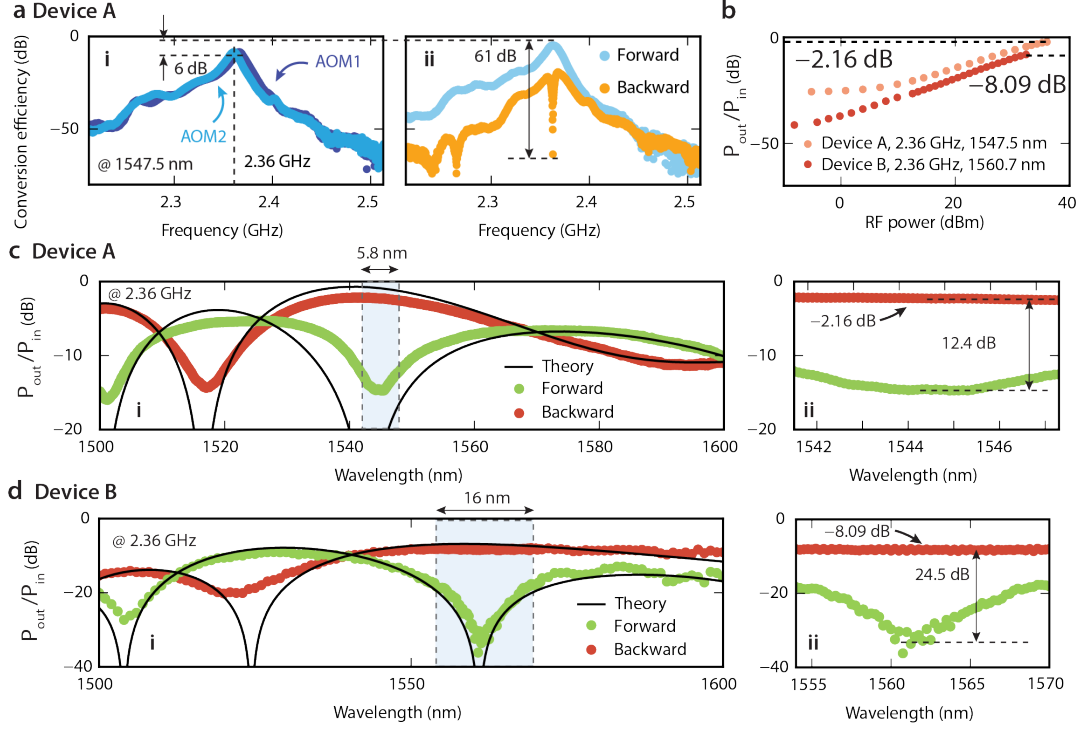


Figure 4.6: **a**, We measure the mode conversion efficiency as a function of the microwave drive frequency while driving only AOM 1 (dark blue in **i**) and driving only AOM 2 (light blue in **i**), which demonstrates almost identical responses. When both AOMs are simultaneously driven (**i**), the forward (backward) operation of our device equates to the sum (difference) of the responses from the two AOMs, leading to the high isolation contrast exhibited by our device. **b**, The insertion loss decreases as we increase the microwave drive power, approaching -2 dB. The maximum RF power is limited by the power handling of the device. **c**, The transmission of Device A operating in forward (red) and backward (green) directions, which demonstrates the 10 dB isolation bandwidth of 5.8 nm (725 GHz, shaded blue window in **i**, enlarged in **ii**). Over the isolation bandwidth, the device has -2.16 dB insertion loss and 12.4 dB peak isolation ratio. Measurements deviate slightly from theoretical predictions (black) due to pump depletion and intermodal crosstalk in the passive photonic components of the device. Note that higher contrast is seen between forward and backward transmission in **(a)** because these frequency-resolved measurements display only the frequency-shifted components of the transmitted light field. **d**, The transmission of Device B operating in forward (red) and backward (green) directions, which demonstrates the 10 dB isolation bandwidth of 16 nm (2 THz, shaded blue window in **i**, enlarged in **ii**). Over the isolation bandwidth, the device has -8.09 dB insertion loss and 28.1 dB peak isolation ratio.

scattered wave components from both AOMs. This frequency is chosen because it ensures equal conversion efficiency from both AOMs, maximizing light cancellation

in the backward direction. To examine the backward operation of the isolator, we inject 1547.5 nm light into Port 4 and measure the optical transmission from Port 1. For the case of backward transmission, the phase-coherent addition of scattered wave components from both AOMs produces destructive interference, corresponding to 61 dB of nonreciprocal contrast at 2.36 GHz drive frequencies. Detuning from the center RF frequency disrupts the phase difference and equal conversion efficiencies which are necessary for perfect destructive interference, leading to a sharp backward dip.

To highlight the potential for both low optical losses and wide isolation bandwidths using this technique, we examine the performance of two devices with the same design, Device A and Device B, having complementary performance characteristics. Device A was released using liquid-phase HF (l-HF) and has higher acousto-optic scattering efficiencies. Since the insertion loss of the isolator is directly determined by acousto-optic conversion efficiency, measurements of isolator insertion loss as a function of microwave drive power (Fig. 4.6b) reveal lower insertion losses at higher microwave drive intensities for both devices. Note that the finite isolator transmission at low microwave drive powers primarily results from intermodal crosstalk in the passive circuit components. As seen from Fig. 4.6b, the higher acousto-optic scattering efficiencies of Device A translate to a lower insertion loss (-2.16 dB) for forward-transmitted light at the maximum microwave drive power. Device B, which was released using vapor-phase HF (v-HF), exhibits lower acousto-optic scattering efficiencies, translating to higher (-8.09 dB) isolator insertion losses (see Fig. 4.6b). However, lower inter-modal crosstalk within the passive circuit elements of Device B translates to higher isolation contrast, and the more optimal waveguide dimensions of Device B also produced improved dispersion balancing, which contributes to wider isolation bandwidths. The distinct AOM performance characteristics of Devices A and B result from the different release methods used.

Next, we analyze the wideband performance of these isolators by sweeping the laser wavelength and measuring the total optical power transmissivity ($P_{\text{out}}/P_{\text{in}}$) while the AOMs are driven at their peak frequencies and maximum RF powers. Transmission measurements in both the forward (red) and backward (green) directions, reveal high-contrast isolation, as is shown in Fig 4.6d for device B (for device A, please see Fig 4.6c). These measurements were performed using direct optical power detection, such that total optical power exiting Port 4 (Port 1) in the forward (backward) direction is compared with the optical power entering Port 1 (Port 4); no spectral filtering is being performed, i.e. crosstalk from mode multiplexer is included. Device B exhibits a peak isolation contrast of 24.5 dB and an insertion loss of -8.09 dB at 1560.7 nm, corresponding to the center of the isolation bandwidth; these spectral measurements reveal a 10 dB isolation bandwidth of 16 nm corresponding to a frequency bandwidth of 2 THz, representing orders of magnitude larger isolation bandwidth than previous nonmagnetic isolators.

This broad isolation bandwidth is enabled by precisely balancing the group velocities of the $\mathbf{E}_+$ and $\mathbf{E}_-$ modes in the dual-ridge waveguide design. The FSR is consistent with a group delay difference of ~ 30 fs; observed variations in the fringe period arise from the distinct group velocity dispersion (GVD) of the two modes within the multimode waveguide. The theoretical isolator performance (black) is based on an analytical model and shows good agreement with the measured isolator response. For Fig. 4.6c, the parameters used in the fitting are $n_{g,+}^d = 3.5729$, $n_{g,-}^d = 3.5721$, $n_{g,+} = 3.805$, $n_{g,-} = 3.885$, $\omega_0 = \frac{2\pi c}{1540.65 \text{ nm}}$, $L_a = 276 \text{ um}$, $L_t = 274 \text{ um}$, $L_d = 7.64 \text{ mm}$, $\beta_{d,+} - \beta_{d,-} = -470 \text{ fs}^2/\text{mm}$, $\phi_{\text{RF1}} - \phi_{\text{RF2}} = -q(L_a + L_t)$. For Fig. 4.6d, the parameters used in the fitting are $n_{g,+}^d = 3.6847$, $n_{g,-}^d = 3.6803$, $n_{g,+} = 3.8628$, $n_{g,-} = 3.9672$, $\omega_0 = \frac{2\pi c}{1558.48 \text{ nm}}$, $L_a = 276 \text{ um}$, $L_t = 274 \text{ um}$, $L_d = 7.64 \text{ mm}$, $\beta_{d,+} - \beta_{d,-} = -372 \text{ fs}^2/\text{mm}$, $\phi_{\text{RF1}} - \phi_{\text{RF2}} = -q(L_a + L_t)$. Due to higher conversion efficiency in Device A, the simplified theory doesn't fit perfectly. For device B, the fitting agrees

better to the experiment.

Notably, this isolation bandwidth can be made even larger with more precise balancing of the dispersion in each arm of the interferometer.

4.3.2 Wavelength tunable wideband isolation based on length balanced singlemode waveguide pair

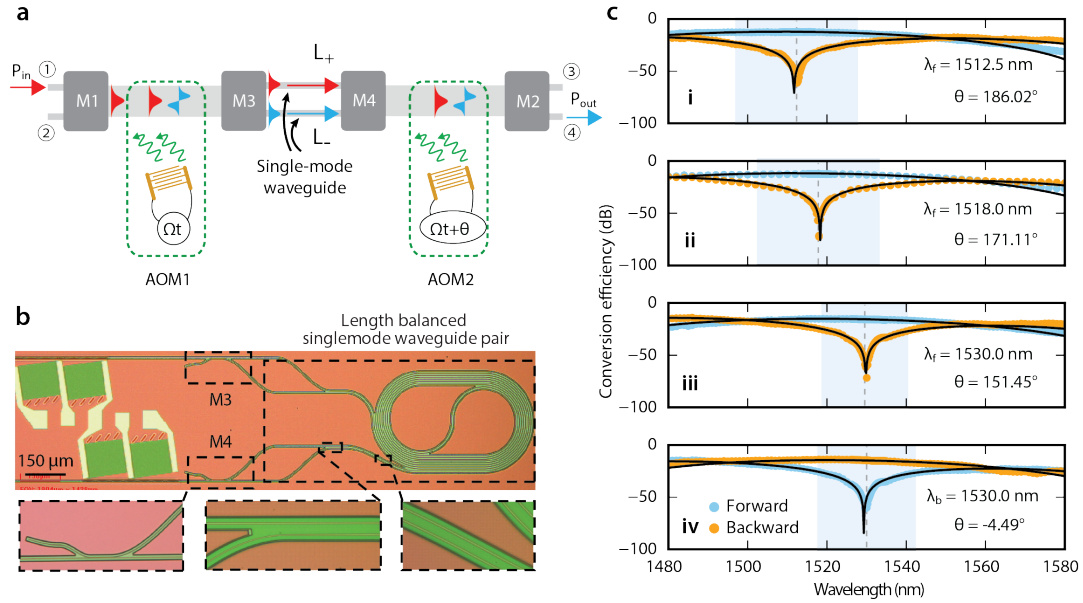


Figure 4.7: **a**, The second dispersion engineering method incorporates two extra mode multiplexers to split the two waveguide modes into two single-mode waveguides. **b**, An optical micrograph of the fabricated device. The two-waveguide delay line is configured in a spiral shape, taking advantage of its robust dispersion properties that can withstand waveguide bends without significant degradation. **c**, Conversion efficiency spectrum with difference phase difference between two AOMs. Our device exhibits wavelength tunability (**i-iii**) and reconfigurability of the isolation direction (**iii** and **iv**), which can be adjusted by controlling the relative phase difference θ between the microwave drives for AOM 1 and AOM 2.

Next, we present a complementary strategy for wideband dispersion engineering, and we use it to demonstrate frequency tunability and rapid reconfigurability of the nonreciprocal response. Through this second approach, in place of a dispersion-engineered multimode waveguide, we instead use two separate single-mode waveguides

(Fig. 4.7a) to precisely match the group delay in each arm, as required to achieve wide-band nonreciprocity. As seen in Fig. 4.7a, waves propagating in modes $\mathbf{E}_+$ and $\mathbf{E}_-$ at the output of AOM 1 are demultiplexed by M3 and transmitted through two identical single-mode waveguides before being recombined by M4 and entering AOM 2. Since both arms of the interferometer are identical in waveguide geometry, this approach produces a constant group-delay difference over all wavelengths, meaning that the fringe period of the interferometer does not change as a function of wavelength.

We examine the nonreciprocal response of the micro-fabricated device seen in Fig. 4.7b. This system exhibits lower acousto-optic scattering efficiencies (-12.29 dB) due to reduced electromechanical couplings with none optimal IDT design in this device. Since the acousto-optic scattering efficiency is not appreciably larger than the intermodal crosstalk, it is not possible to demonstrate optical isolation using direct power detection, as we did in Fig. 4.6 for the design based on dispersion-engineered multimode waveguide. Instead, we use optical heterodyne measurements to analyze the constructive and destructive interference between the acousto-optically scattered tones and illustrate the the unique advantages of this group delay balancing approach. Fig.4.7c.i shows the frequency-resolved transmission measurements (i.e. crosstalk removed) in forward and backward directions while both AOM 1 and AOM 2 are electrically driven. In the forward direction (blue), the frequency-shifted components of the transmitted wave produce constructive interference, with maximum of conversion efficiency at 1515.4 nm. In the backward direction (orange), we observe cancellation of the frequency-shifted wave components, corresponding to 48.7 dB of nonreciprocal contrast at 1512.5 nm. The wavelength-dependent transmission in both forward and backward directions show excellent agreement with the predicted nonreciprocal response (black), revealing a 10 dB nonreciprocal bandwidth of 31 nm, highlighted in Fig. 4.7c.i. Note that the nonreciprocal contrast is much larger than those of Fig. 4.6 since these frequency-resolved transmission measurements artificially eliminate inter-

modal crosstalk associated with the passive components within the system.

This wide nonreciprocal bandwidth is enabled by the precise balancing group delay produced in each arm of the interferometer; the fitted FSR of 83 nm corresponds to a 31-femtosecond difference in group delay between the optical paths. To achieve this, it was necessary to use slightly different waveguide lengths ($L_+ - L_- = 10.6 \text{ } \mu\text{m}$) in upper and lower arms of the interferometer to compensate for dissimilarities in the group delay experienced by modes $\mathbf{E}_+$ and $\mathbf{E}_-$ as they traverse the acousto-optic modulators and mode multiplexers (see Supplementary Information III for details). In principle, the nonreciprocal bandwidths of this system can be dramatically enhanced by more precisely adjusting the relative lengths of the upper and lower waveguides within the interferometer to improve further the group delay balancing in the system.

Furthermore, the nonreciprocal passband and direction are readily reconfigured by controlling the relative phase of the microwave tones driving AOM 1 and AOM 2. Since the acousto-optically scattered wave components inherit the phase of the microwave drive tone, we can shift the interferometer fringes as well as the passband center frequency by tuning the relative phase of the microwave tones that are used to drive AOM 1 and AOM 2. Figures 4.7c.i-iii show a sequence of frequency-resolved transmission measurements as the microwave phase (θ) at the input of AOM 1 is varied by 35 degrees, tuning the nonreciprocal passband by 18 nm. The tuning range of the passband is currently limited by wavelength dependence of AOM conversion efficiency (Supplementary Information Fig. S1). Moreover, by implementing a 180-degree change in the phase of the microwave drive, as seen in Fig. 4.7c.iv, we see that the direction of nonreciprocal light transmission becomes reversed, enabling rapid reconfigurability of the direction of isolation.

These two methods for extending the isolator bandwidth— using a single dispersion-balanced multimode waveguide or length-balanced singlemode waveguide pair— have distinct advantages. While we have demonstrated low losses and wideband optical iso-

lation using the dispersion-engineered multimode waveguide, control of the waveguide geometry with nanometer-scale precision is necessary, which is not always practical. We have shown that comparable nonreciprocal bandwidths can be achieved using two identical waveguides to ensure that the group delay produced in each arm is well-matched. This approach eliminates the need for tight dimensional control of the waveguide geometry; both waveguides need only have identical dimensions, a condition easily met with most fabrication processes. However, this second approach places more stringent requirements on mode multiplexer performance to achieve high contrast isolation. Nevertheless, with further design optimization, either approach can be used to produce high-performance wideband isolators in any number of photonic platforms (Fig. 4.8). Compared with traditional traveling acoustic wave modulation approach [27], our results achieve 30 times larger bandwidth. Moreover, compared resonator based approaches, our results show more than 1000 times bandwidth increase.

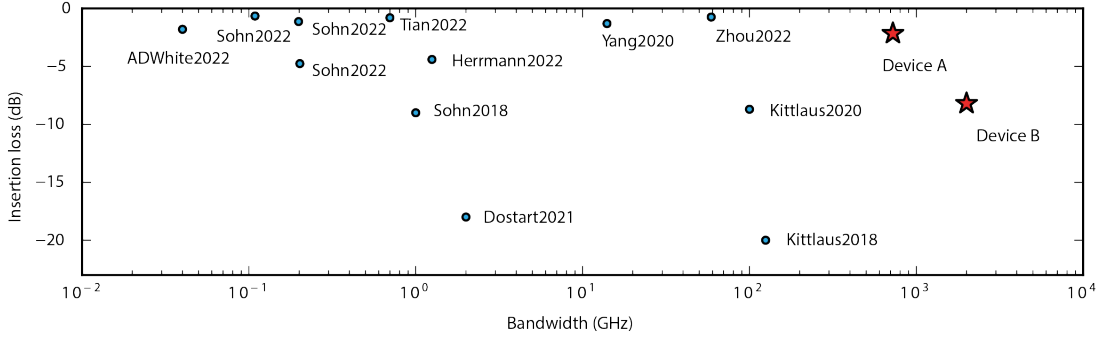


Figure 4.8: Review of on-chip non-magnetic optical isolators. Collection of the previous results shown in [22–27, 82–84].

4.4 Discussion and conclusion

In conclusion, we have demonstrated a new approach for wideband nonmagnetic optical isolation that combines nonreciprocal light scattering processes with dispersion

engineering to realize a bandwidth-scalable nonmagnetic isolator. Using this approach, we have demonstrated high isolation contrast over bandwidths as high as 2 THz, representing a major advance for nonmagnetic isolators in integrated photonics. These results correspond to a 30-fold increase in bandwidth over waveguide-based devices [26,27] and a over 1000-fold increase relative to resonator-based devices [22–24,84]. Moreover, with more precise group delay balancing, isolation bandwidths exceeding 10 THz can be achieved, extending the capabilities of such integrated non-reciprocal technologies.

Further improvements in the AOM performance will translate to lower isolator insertion losses and greatly reduced power consumption, which is necessary for practical integrated technologies. The insertion losses of these isolators are currently limited by the AOM scattering efficiency. Hence, a modest increase in AOM scattering efficiency would eliminate this source of excess loss, reducing the insertion losses to ~ 0.8 dB levels (limited by waveguide loss). Additionally, with further refinement of the electromechanical transducers that drive the AOMs, the microwave power necessary to power this isolator can be reduced by several orders of magnitude. Simulations reveal that ~ 400 μ W of acoustic power guided within the optomechanical waveguide is sufficient to produce complete acousto-optic conversion within the AOMs. Hence, with further optimization of the electromechanical transducer design it should be possible to reduce the power consumption of the isolator to milliwatt levels.

Looking beyond this demonstration, these concepts for creating and controlling wideband nonreciprocity are broadly applicable to other photonic platforms. While we used acousto-optic scattering to make this demonstration, other intermodal scattering mechanism, such as electric-optical modulation [19], can in principle be used to achieve the same functionality, opening the door to wideband isolation using a broader range of interactions in a variety of materials platforms. Rapid electrical reconfigurability and tunability of the isolator passband also offer intriguing possibilities

for extending the bandwidth and functionalities of photonic systems, with potential impacts in sensing, communications, computing, and frequency metrology.

Chapter 5

Conclusion

5.1 Summary

In this dissertation, we have presented key advancements in silicon photonics, with a focus on integrated laser stabilization, nonmagnetic optical isolation, and their applications in photonic integrated circuits (PICs). By addressing fundamental challenges in laser coherence, optical feedback control, and isolation, this work has contributed to the development of highly stable, low-noise, and CMOS-compatible photonic systems with the potential to impact optical communication, precision sensing, and quantum technologies.

A central achievement of this dissertation is the co-integration of high-performance μ FP cavities with PICs to enable superior laser stabilization. While bulk optical reference cavities offer excellent frequency stability, their large size and need for complex vacuum enclosures hinder integration into chip-scale systems. In contrast, this dissertation demonstrates a novel approach to integrating μ FP cavities with silicon photonics, employing wafer-scale fabrication techniques to create compact, vacuum-sealed μ FP cavities with ultrahigh quality factors ($> 10^9$) and record-low phase noise. A key enabler of this integration is the reflection transformation circuit, which not

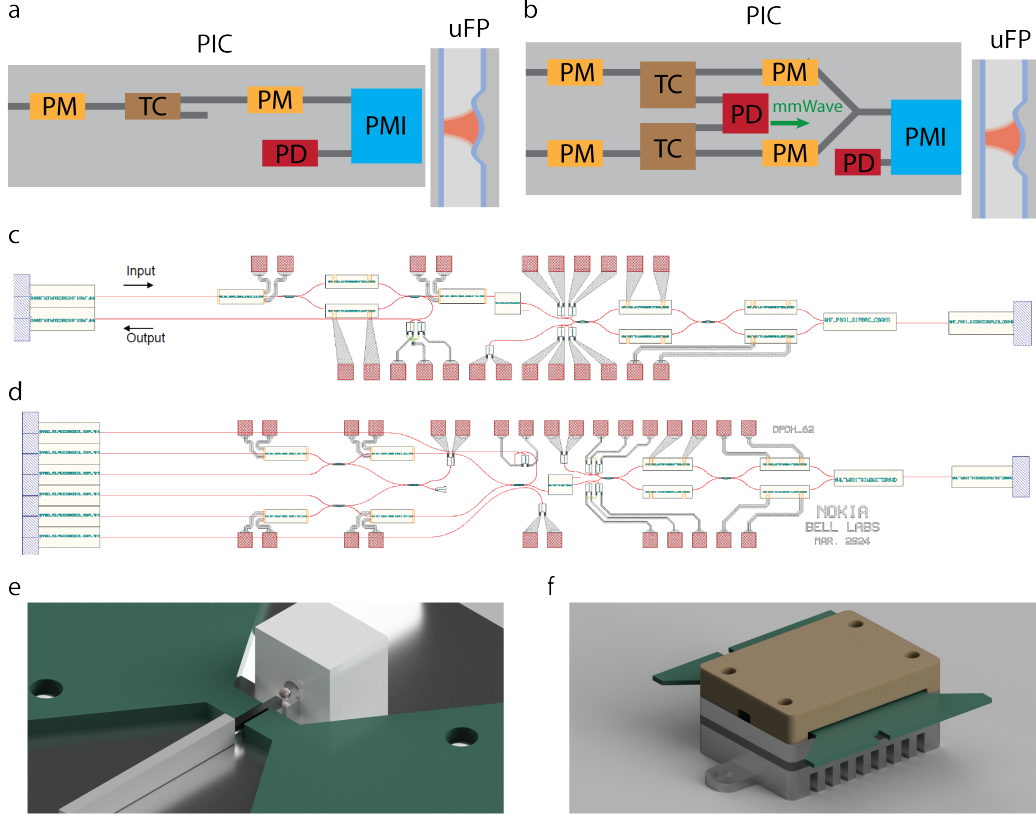


Figure 5.1: **a** On-chip PDH locking circuit design. PM: phase modulator, TC: tunable coupler, PD: photo detector, PMI: poor man’s isolator we discussed in Sec 3.3.1. **b** Dual PDH locking for mm-wave generation. **c** Detailed GDS design of on-chip PDH locking chip. **d** Detailed GDS design of on-chip dual-PDH locking chip. **e** Zoom-in of the sub-mount and package for dual-PDH chip. **f** Zoom-out of the package for dual-PDH chip.

only mitigates the mirror-like reflection inherent to FP cavities but also provides precise optical feedback control, eliminating the need for bulky, non-CMOS-compatible optical circulators. This innovation ensures seamless incorporation of μ FP cavities into silicon photonic platforms while preserving the isolation required for stable laser operation.

By implementing self-injection locking (SIL) within integrated photonic platforms, we have demonstrated record-breaking phase noise suppression and frequency stability improvements. Specifically, the SIL-integrated laser achieved a single-sideband phase noise of -97 dBc/Hz at 10 kHz offset and a fractional frequency stability of

5×10^{-13} at 10 ms, representing a significant leap in integrated laser performance. Additionally, a co-integrated μ FP cavity was demonstrated to enable on-chip electrical feedback, effectively suppressing back-reflections and redirecting the reflected signal for detection. With a suppression ratio exceeding 10 dB, this advancement facilitates the seamless implementation of on-chip Pound-Drever-Hall (PDH) locking, making high-performance laser stabilization more accessible within integrated photonics.

Beyond laser stabilization, this dissertation introduces a new approach to achieve on-chip nonmagnetic optical isolation, leveraging acousto-optic interactions. Conventional optical isolators rely on magneto-optic materials such as yttrium iron garnet, which are incompatible with CMOS fabrication processes, posing a major challenge for full integration. To address this, this work develops a nonmagnetic isolation scheme based on dispersion-balanced multimode interferometers and engineered intermodal acousto-optic scattering processes. By carefully designing the optical waveguide delay and optimizing acousto-optic modulation, this approach achieves broadband isolation as high as 2 THz at 1550 nm, with insertion loss as low as 2 dB. These results establish a scalable, on-chip optical protection solution, eliminating the reliance on magneto-optic materials and enabling fully CMOS-compatible optical isolation.

The advancements demonstrated in this dissertation have far-reaching implications for the future of integrated photonics. The ability to co-package high-quality optical reference cavities with silicon photonic circuits unlocks new possibilities for ultra-low-noise microwave generation, high-precision optical clocks, and quantum photonics. At the same time, the development of wideband nonmagnetic optical isolators addresses a critical need for integrated laser protection, ensuring the stable operation of chip-scale laser sources in applications ranging from optical communication networks to LIDAR and precision metrology. Together, these innovations establish a foundation for the next generation of high-performance silicon photonic

systems, bridging the gap between laboratory demonstrations and scalable, real-world applications.

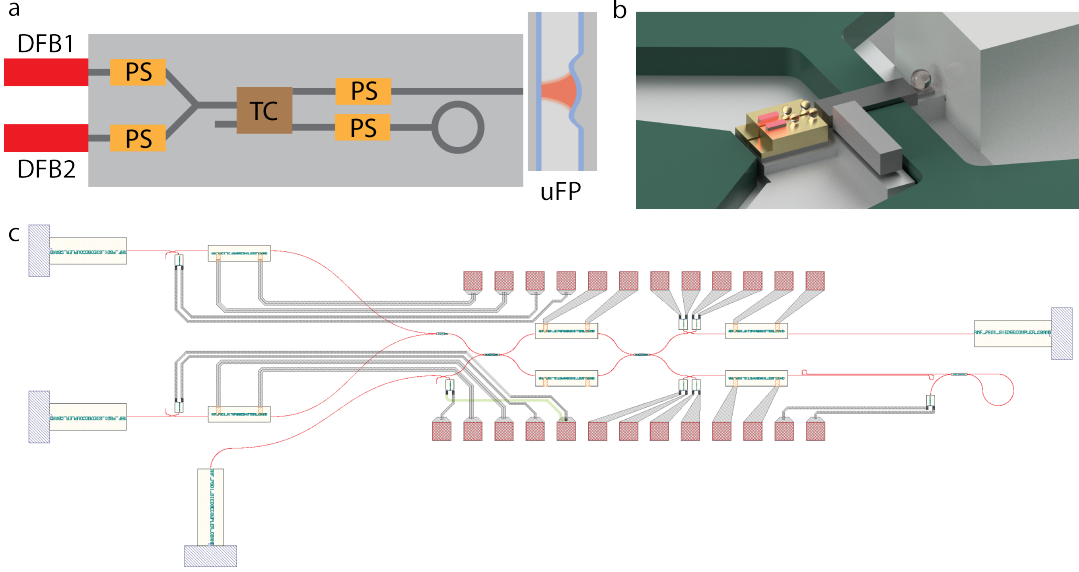


Figure 5.2: **a** On-chip dual-SIL circuit design. PS: phase shifter, TC: tunable coupler, PD: photo detector. **b** Zoom-in of the package design, with two DFB lasers sitting on two separate TECs. **c** Detailed GDS design of on-chip dual-SIL chip.

5.2 On-going projects and outlook

Building on these achievements, we designed photonic integrated circuits and fabricated them through a commercial silicon photonic foundry to further explore the applications of on-chip ultra-stable systems. As a first step, leveraging commercially available on-chip modulators and detectors, we developed an integrated PDH locking circuit (Fig.5.1(a,c)) to demonstrate the first fully integrated PDH-locked laser, as previously proposed in literature [37, 64]. Building on this foundation, we extended the approach to dual-PDH locking (Fig.5.1(b,d)), where two independent lasers are locked to two different longitudinal modes of the FP cavity, separated by one free spectral range. This configuration allows the two ultra-stable lasers to heterodyne, generating an ultra-stable millimeter-wave (mm-wave) signal.

Similarly, we implemented dual self-injection locking (SIL) by locking two distributed feedback (DFB) lasers to the same μ FP cavity at two distinct longitudinal modes (Fig. 5.2(a-c)). This dual-SIL approach eliminates the need for complex electronic control, while simultaneously offering a higher noise suppression bandwidth. These advancements pave the way for fully integrated ultra-stable optical and microwave signal generation, with significant implications for applications in high-precision sensing, coherent communication, and frequency metrology.

Furthermore, we designed a circuit to address a key limitation of SIL lasers—the lack of long-term stability and their inability to correct for frequency drift. When the DFB laser frequency drifts beyond the SIL locking range, the laser loses its locked state, disrupting stable operation. To overcome this issue, we leverage a key characteristic of the SIL interface circuit—its output port exhibits a reflection dip response, making it well-suited for use as a frequency discriminator [44]. By introducing a beat signal with a reference arm (Fig. 5.3(a)), a balanced photodetector (BPD) can determine whether the DFB laser is blue- or red-detuned relative to the μ FP cavity. Using this information, an electrical feedback loop can be implemented to automatically correct frequency drift, ensuring continuous SIL operation. This design significantly enhances the long-term stability of SIL lasers, making them robust and reliable for industrial applications, particularly in high-precision sensing, frequency metrology, and optical communication systems where stable, ultra-narrow linewidth lasers are essential.

Looking ahead, the impact of these research extend beyond the laboratory. The integrated platforms and methodologies developed here are directly applicable to high-speed optical communications, precision metrology, quantum information processing, and emerging photonic microwave systems. Future work may explore further miniaturization of integrated components, the use of alternative materials for even lower loss, and the integration of these photonic circuits with electronic and quantum

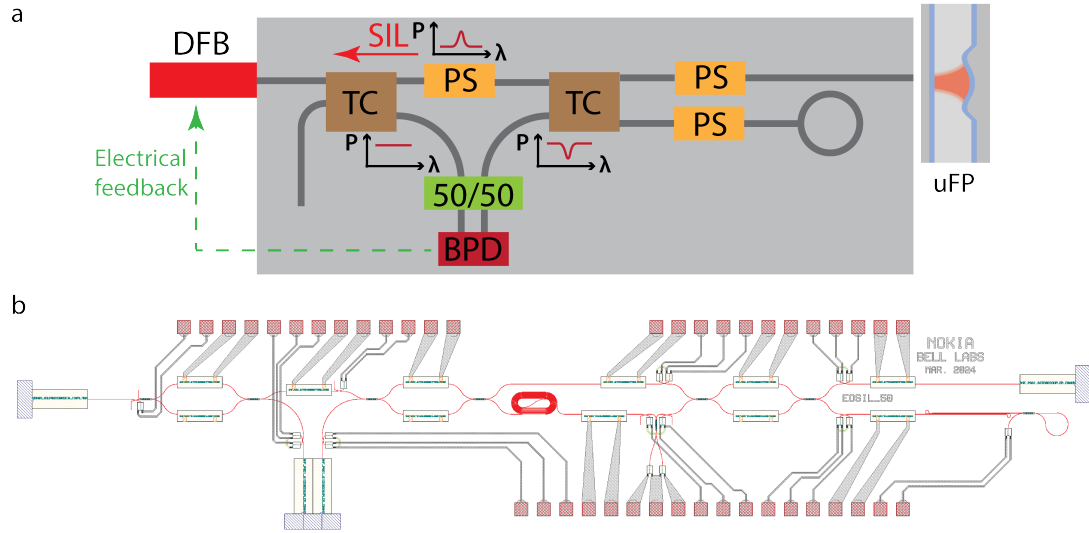


Figure 5.3: **a** Schematic of hybrid electro-SIL laser. BPD: balanced photo detector **b** Detailed GDS design of hybrid electro-SIL laser chip.

systems to create hybrid devices with unprecedented functionality.

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