



Dispersion Engineering of Thin-Film Lithium Niobate Waveguides for Quantum Applications

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Summary

Scalable and efficient devices are crucial for the miniaturization of integrated quantum optics experiments. A new frontier material in this research area is thin-film lithium niobate (TFLN). However, a systematic investigation of the feasibility of the material for this research field is not fully completed. Thus, this thesis systematically assesses TFLN as a material for integrated quantum optics by harnessing its advantages for dispersion engineered quantum devices.

We investigate propagation losses in TFLN waveguides as the basic building block for integrated circuits showing state of the art losses. We demonstrate a Hong-Ou-Mandel interference in a TFLN directional coupler as a simple quantum optical phenomenon. Since photon sources are another important building block, we simulate and characterize different dispersion engineered parametric down-conversion sources in-depth to harness the advantages of highly confined waveguides in lithium niobate. The full potential of TFLN takes shape with the realization of a highly efficient dispersion engineered quantum pulse gate that outperforms standard devices by three orders of magnitude.

Using this systematic approach, we demonstrate the applicability of TFLN for integrated quantum optics and establish an infrastructure for the design and characterization of quantum devices laying the foundation for the realization of highly sophisticated quantum circuits on TFLN.

Zusammenfassung

Skalierbare und effiziente Bauelemente sind wesentlich für die Miniaturisierung integrierter quantenoptischer Experimente. Ein aussichtsreiches neues Material in diesem Forschungsfeld ist Dünnschicht Lithiumniobate (TFLN). Jedoch ist eine systematische Untersuchung des Materials für dieses Forschungsfeld noch ausstehend. Daher untersucht diese Arbeit TFLN systematisch als Material für integrierte Quantenoptik, indem die Vorteile für dispersionsoptimierte Quantenbauelemente genutzt werden.

Wir untersuchen die Propagationsverluste in TFLN-Wellenleitern als ein Baustein für integrierte Schaltkreise und erreichen State of the Art Verluste. Wir demonstrieren Hong-Ou-Mandel-Interferenz in einem TFLN-Richtkoppler als einfaches quantenoptisches Phänomen. Da Photonenquellen ein weiterer wichtiger Baustein sind, simulieren und charakterisieren wir verschiedene dispersionsoptimierte parametrische Fluoreszenzquellen detailliert, um den Vorteil von Wellenleiter mit einem hohen Indexkontrast in Lithiumniobat zu nutzen. Das ganze Potenzial von TFLN zeigt sich mit der Realisierung eines hocheffizienten, dispersionsoptimierten Quantenpulsgatters, das Standardbauelemente um drei Größenordnungen übertrifft.

Mit diesem systematischen Ansatz zeigen wir die Anwendbarkeit von TFLN für integrierte Quantenoptik und schaffen eine Infrastruktur für das Design und die Charakterisierung von Quantenbauelementen. Dies bildet die Grundlage für die Realisierung hochkomplexer Quantenschaltungen in TFLN.

1

Introduction

We are currently living in the era of the second quantum revolution [1, 2]. In contrast to the first quantum revolution in which the main focus was to understand quantum phenomena, an effort is now being made to actively harness quantum states to develop new technologies. These technologies include quantum computers, quantum communication schemes as well as quantum metrology [3]. To implement any of those technologies, a suitable information carrier is needed, a promising choice are photons. Photons have the advantage that they do not have a photon-photon interaction and interact only weakly with the environment. This leads to a low decoherence and a low-loss transmission via the already existing fiber network. This makes photons an ideal carrier for quantum information.

Most of the newly developed technologies in photonics are based on nonlinear optical processes to generate single photons or convert frequencies. The standard processes to generate entangled photons are three-wave mixing processes such as parametric down-conversion or four-wave mixing processes [4]. To harness the full potential of nonlinear devices processes based on a second-order nonlinearity are favorable due to the enhanced efficiency compared to third-order nonlinear processes. One important material which exhibits a second-order nonlinearity is lithium niobate.

Since the 1960s lithium niobate is a widely used material in photonics [5]. In the beginning, bulk lithium niobate was used to carry out fundamental optical experiments. These experimental setups filled complete optical tables. To miniaturize and stabilize the experiments, soon waveguides were fabricated in lithium niobate either by proton exchange or titanium indiffusion [6, 7]. With these waveguides, experiments were miniaturized and single experiments such as an on-chip Hong-Ou-Mandel interference were demonstrated [8].

However, the indiffused waveguides in lithium niobate exhibit only a weak index contrast leading to large bending radii. Consequently, complex experiments consisting of multiple functionalities cannot be integrated on a single wafer. This means a further miniaturization to realize scalable photonic circuits on weakly confined waveguides in lithium niobate is not feasible compared to high index contrast waveguides such as silicon photonics.

Thus, to achieve smaller footprint devices, we have to use a different material platform which combines a high integration density with a second order nonlinearity. This material platform is called thin-film lithium niobate (TFLN) and was first introduced by Levy et al. in 1998 [9]. First devices based on the material platform followed shortly after [10, 11, 12, 13, 14]. This material platform was also investigated already in 2009 in Paderborn by Hui Hu and coworkers [15] who later founded the company *NANOLN* which is currently one of the main suppliers of TFLN wafers.

TFLN consists of a thin layer of lithium niobate which is bonded onto an insulator layer

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which is connected to a handle substrate. One approach is to guide the light in etched waveguides in the lithium niobate thin-film resulting in the same optical properties as lithium niobate. At the same time, the geometry leads to a high refractive index contrast between lithium niobate and silicon dioxide/air allowing an up to two orders of magnitude higher index contrast change compared to standard lithium niobate waveguides [16]. This leads to a bending radius in the micrometer range. Consequently, small footprint devices on chip leading to the possibility of designing multi-functional circuits can be realized. Moreover, due to the high mode confinement, TFLN creates the opportunity of dispersion engineering since the dispersion does not only depend on the material dispersion as it is the case for low index contrast waveguides, but it strongly depends on the waveguide geometry [17]. This opens up the possibility to design novel devices with tailored properties compared to standard lithium niobate.

At the beginning of TFLN, the etched high index contrast waveguides had losses of 6.3 dB/cm [15] making them not suited for quantum applications. The start of TFLN as a material for integrated quantum optics was in 2017 when Zhang et al. measured losses of only 0.027 dB/cm [18]. Shortly after this, the first PDC source in TFLN was realized in 2019 by Chen et al. [19]. Showing the generation of single photons in TFLN waveguides was an important step for TFLN as a material platform for integrated quantum optics. Since then many devices have been demonstrated in this material platform such as electro-optic modulators [20, 21, 22, 23, 24] or frequency converters [25, 26, 27, 16].

However, at the start of this thesis, only a few quantum devices had been shown in TFLN such as PDC sources [19, 28, 29, 30]. Thus, what was missing at that point was a systematic investigation of the material platform with the focus on quantum applications.

Therefore, in this thesis, we systematically investigate TFLN as a material for integrated quantum optics, we show new functionalities of TFLN devices and establish an infrastructure for the design and characterization of quantum devices based on TFLN.

Structure of the Thesis

This thesis explores the benefits, in particular dispersion engineering, of TFLN waveguides for quantum applications. In chapter 2, we set the scene by explaining the theoretical concepts. First, we start with explaining the working principle of waveguides as the basic device of integrated quantum optics. Next, we explore integrated devices and discuss directional coupler as well as resonators where we explain two common methods to measure propagation losses. With this we move on to nonlinear phenomena and explain three-wave mixing processes and phase matching. As a next step, we introduce the material platform thin-film lithium niobate in which we will use to integrate the discussed devices and effects. This chapter is concluded by delving into the topic of quantum optics where we elaborate on parametric down-conversion, Hong-Ou-Mandel interference and the quantum pulse gate. The next chapter deals with how to measure propagation losses in TFLN waveguides. Here, we first give an overview of measured propagation losses in TFLN waveguides. Then, we display the experimental setup and we discuss the design of the samples. Lastly, we show and compare the measurement results.

In the following chapter, we dive into the quantum part of this thesis. First, we discuss the design and characterization of directional couplers. After showing the newly built experimental setup, we discuss the measured HOM dip.

In chapter 5 the topic of dispersion engineering is introduced. We first explain the principle of dispersion engineering in TFLN waveguides. Then, we will discuss the design and

experimental characterization of a highly non-degenerate PDC in TFLN. Afterwards, we consider the topic of fabrication tolerances on the phase matching spectrum and examine method to counteract those. Lastly, we design sources with specific group velocity matching conditions.

In the sixth chapter, we demonstrate the first QPG in TFLN. Here, we will first elaborate on the dispersion engineering of the devices. Then, we will explain the setup and map the phase matching function. Lastly, we benchmark our device by evaluating the temporal mode selectivity as well as the conversion efficiency.

The final chapter provides a thorough conclusion based on the discussed results and gives an outlook on possible future research.

2

Foundations and Theoretical Concepts

This thesis aims to explore dispersion engineering for quantum optical applications in TFLN waveguides. Thus, to start we lay some theoretical foundation to explain the following work. In this chapter, we first discuss the theory of waveguides as they are the fundamental building block of this thesis. Afterwards, we will move on to more complex integrated structures such as directional couplers and resonators. Then, we will delve into the field of nonlinear optics and gain insight into important nonlinear processes and principles such as three-wave mixing processes and phase matching. Next, we will introduce the used material platform and describe the simulation and fabrication of waveguides. Lastly, we will use this knowledge to explore the principles of quantum optical processes in particular parametric down-conversion, Hong-Ou-Mandel interference and the quantum pulse gate.

2.1 Theory of Waveguides

The key component of an integrated optical circuit is a waveguide to guide and route the light. Thus, we will start this thesis with reviewing the working principle and key features of waveguides. To do so, we will start with a mathematical description of a waveguide. Furthermore, we will use a quantum mechanical analogon to understand the properties of a waveguide and we discuss 2D waveguide features.

2.1.1 Mathematical Description of Waveguides

The basic functionality of a waveguide can be understood by looking at a planar homogeneous waveguide. Such a waveguide consists of three different layers called the substrate with the refractive index n_s , followed by the film with n_f and the cover n_c . The light is guided in the film due to total internal reflection [31].

To describe a waveguide mathematically, we start with Maxwell's equations and from this we will get to the vectorial wave equations. Based on these, we will discuss for a homogeneous planar waveguide the field distribution of modes inside the waveguide.

First, we write down Maxwell's equations for a dielectric, non-magnetic and non-conducting

Chapter 2. Foundations and Theoretical Concepts

material [31]:

$$\nabla \times \vec{E} = -\frac{\partial \vec{B}}{\partial t} \quad (2.1)$$

$$\nabla \times \vec{B} = \mu_0 \frac{\partial \vec{D}}{\partial t} \quad (2.2)$$

$$\nabla \cdot \vec{D} = 0 \quad (2.3)$$

$$\nabla \cdot \vec{B} = 0. \quad (2.4)$$

Here, $\vec{E}$, $\vec{B}$ and $\vec{D}$ describe the electric field, the magnetic induction and the electric displacement, respectively and μ_0 describes the magnetic permeability. We assume monochromatic waves in our waveguide with the angular frequency ω of the form:

$$\vec{E}(x, y, z, t) = \vec{E}(x, y, z) \cdot e^{i\omega t} \quad (2.5)$$

$$\vec{B}(x, y, z, t) = \vec{B}(x, y, z) \cdot e^{i\omega t}. \quad (2.6)$$

Inserting these fields into the first two Maxwell equations, we can re-write them as [31]:

$$\nabla \times \vec{E} = -i\omega \vec{B} \quad (2.7)$$

$$\nabla \times \vec{B} = \mu_0 i\omega \vec{D}. \quad (2.8)$$

To get to the vectorial wave equation, we apply a curl operator and make use of the following mathematical identity $\nabla \times (\nabla \times \vec{E}) = \nabla (\nabla \cdot \vec{E}) - \Delta \vec{E}$ [32]. Applying this to Equation 2.7 leads us to:

$$\nabla \times (\nabla \times \vec{E}) = i\omega \nabla \times \vec{B} \quad (2.9)$$

$$\nabla (\nabla \cdot \vec{E}) - \Delta \vec{E} = \omega^2 \mu_0 \vec{D}. \quad (2.10)$$

To simplify this equation even more, we can assume that our waveguide is made out of an isotropic material which means that the permittivity is a scalar and not a tensor. If this is the case, we have the following relation between the electric field and the electric displacement [31]:

$$\vec{D} = \epsilon_0 \epsilon \vec{E}. \quad (2.11)$$

Here, ϵ_0 is the vacuum permittivity and ϵ the dielectric permittivity. If we insert this relation into Equation 2.10 we get:

$$\nabla \cdot \vec{E} = -\frac{\nabla \epsilon}{\epsilon} \cdot \vec{E}. \quad (2.12)$$

As a last step, we combine Equation 2.12 and Equation 2.10 to get to the vectorial wave equation [33]:

$$-\nabla \left(\frac{\nabla \epsilon}{\epsilon} \cdot \vec{E} \right) - \Delta \vec{E} = \omega^2 \mu_0 \epsilon_0 n^2 \vec{E}. \quad (2.13)$$

Here, we also used the relation between the refractive index n and dielectric permittivity ϵ : $n^2 = \epsilon$. Applying the same steps to Equation 2.2 leads to the vectorial wave equation for the magnetic field:

$$\Delta \vec{B} + \left(\frac{\Delta \epsilon}{\epsilon} \times (\Delta \times \vec{B}) \right) + \mu_0 \epsilon_0 n^2 \omega^2 \vec{B} = 0. \quad (2.14)$$

2.1. Theory of Waveguides

With these equations the field distribution for different waveguide geometries with a given refractive index distribution can be calculated. For most waveguides, numerical methods are necessary to solve the equations as we will later do for our TFLN waveguides. However, we can apply these equations to our basic model of a homogeneous planar waveguide and calculate the field distribution by hand. The results can also be applied to more complex waveguide structures. Therefore, in the following we will solve the vectorial wave equation for a planar homogeneous waveguide.

For this, we assume that we have a wave that propagates along the z -axis. The refractive index of our material is constant in the y -direction and thus changes only in the x -direction. This yields that all derivatives in y -direction are zero. A wave propagating in z -direction can be written as [33]:

$$\vec{E}(x, y, z, t) = \vec{E}(x, y) \cdot e^{i(\beta z - \omega t)} \quad (2.15)$$

$$\vec{B}(x, y, z, t) = \vec{B}(x, y) \cdot e^{i(\beta z - \omega t)}. \quad (2.16)$$

Here, β is the propagation constant which describes the vacuum wave number component k in propagation direction [31]:

$$\beta_i = \frac{\omega}{v_p} = k \cdot n_f \cdot \sin(\Theta_i) = k \cdot n_{\text{eff},i}. \quad (2.17)$$

Here, v_p is the phase velocity and n_f the refractive index of the film. The propagation constant can be defined for each guided mode of the waveguide indicated by the index i . The variable n_{eff} is called the effective refractive index and describes the refractive index "seen" by the mode propagating in the waveguide. Inserting these fields into Equation 2.1 and Equation 2.2 yields:

$$\left. \begin{aligned} -i\beta E_y &= -i\omega B_x \\ i\beta E_x - \frac{\delta E_z}{\delta x} &= i\beta B_y \\ \frac{\delta E_y}{\delta x} &= i\omega B_z \end{aligned} \right\} TE$$

$$\left. \begin{aligned} -i\beta B_y &= -i\omega \mu_0 \epsilon \epsilon_0 E_x \\ i\beta B_x - \frac{\delta B_z}{\delta x} &= -i\omega \mu_0 \epsilon \epsilon_0 E_y \\ \frac{\delta B_y}{\delta x} &= -i\omega \mu_0 \epsilon \epsilon_0 E_z \end{aligned} \right\} TM.$$

These two sets of equations are independent from each other and only certain field component combinations solve these. The first combination can be seen if we set $E_x = 0$ since this yields that $E_z = 0$ and $B_y = 0$. This combination of fields is called the transverse electric (TE) modes. The other combination is called transverse magnetic (TM) modes, since by setting $B_x = 0$ only E_x , E_z and B_y are non zero. As the names indicate for the TE modes the electric field is perpendicular to the propagation direction and for the TM modes the magnetic field is perpendicular to the propagation direction. [33]

We will now use the relation we get for the TE modes to solve Equation 2.13. For this, we first assume that we have a constant refractive index in each layer of the waveguide, which means that $\nabla \epsilon = 0$. From the second term of this equation we get $\Delta \vec{E} = \frac{\partial^2 E_y}{\partial x^2} - \beta E_y$ if we take into account that all derivatives in the y -direction are vanishing. This gets us to the Helmholtz equation [33]:

$$\frac{\partial^2 E_y}{\partial x^2} + (n^2(x) \cdot k - \beta^2) E_y = 0. \quad (2.18)$$

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The equation has two solutions depending on the sign of $n^2(x) \cdot k - \beta^2$. For a positive sign the solution is given by [33]:

$$E_y = A \cdot \cos(k_x \cdot x) + B \cdot \sin(k_x \cdot x). \quad (2.19)$$

For a negative sign, we get [33]:

$$E_y = C \cdot e^{(\beta^2 - n^2 \cdot k^2) \cdot x} + D \cdot e^{-(\beta^2 - n^2 \cdot k^2) \cdot x}. \quad (2.20)$$

From these solutions we obtain three different types of modes in a waveguide. If we have a radiation mode, the sign of β is positive in all three layers of the waveguide. This means we have oscillating fields in all layers. In the case of a substrate mode, we have a negative β in the cover and positive in substrate and film. Thus, we have an oscillating mode in the substrate and film and an exponentially decaying wave in the cover. If we have a guided mode, we have a positive β in the film and a negative one in substrate and cover, leading to exponentially decaying fields in the substrate and cover, but an oscillating field in the film. [33, 31]

This analysis can also be performed in a similar manner for the TM modes which yields similar results.

The mathematical approach is one way to understand modes in a waveguide. However, we can also use a quantum mechanical approach to describe how modes occur in a waveguide with which we can also explain other phenomena such as losses in a waveguide in an easier manner. Therefore, in the next part, we will discuss how we can use the principle of a finite potential well to understand the behavior of waveguides.

Finite Potential Well

A waveguide can also be seen as a finite potential well since the scalar wave equation (see Equation 2.18) is formally identical with the Schrödinger equation. Therefore, the well of the finite potential well is composed of the refractive index profile of the waveguide and the solution of the well are the modes of the waveguide. We can attribute the bounded solutions of the well to guided modes of the waveguide and unbounded to substrate modes.

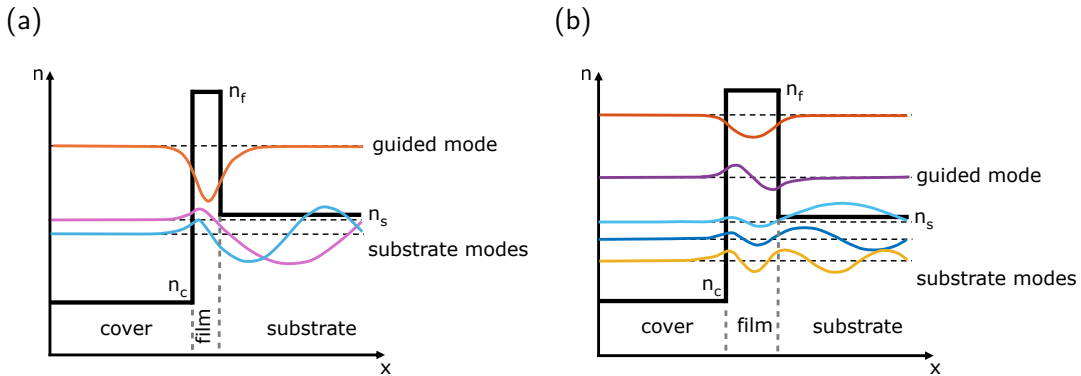


Fig. 2.1: Finite potential well analogon of a waveguide. (a) The panel shows a single-mode waveguide as a potential well where only one guided mode forms. (b) The right panel shows a multi-mode waveguide in which two modes are guided in the waveguide.

2.1. Theory of Waveguides

Depending on the size of the well, it can have one solution which corresponds to a single-mode waveguide (see Figure 2.1 (a)). If the well is larger, the waveguide is multi-mode since more bounded solutions are possible as it can be seen in Figure 2.1 (b). For a guided mode we can see the exponential decay in the substrate and film as calculated with the wave equation. This exponentially decaying field is called the evanescent field. The evanescent field decays faster the higher the index contrast. In the film, we see the oscillating fields as discussed above for substrate modes.

To conclude, in a homogeneous planar waveguide three different types of modes can form. The first one is a radiation modes which oscillates in all three layers of the waveguide [34]. This means that the mode is not guided and experiences losses since the energy radiates away into the substrate and cover. A substrate mode is supported by the waveguide, but since it is oscillating at the film and substrate interface, the mode loses continuously energy such that it only propagates for a short distance [34]. Only a guided mode which is confined in the film will be used in this thesis since this mode can propagate through a waveguide while experiencing less losses compared to the other two mode types. The kind of losses even a guided mode can experience will be discussed next.

2.1.2 Origin of Losses in Waveguides

The fabrication of waveguides often leads to rough sidewalls as it can be seen in Figure 2.2. In this scanning electron microscope (SEM) image two etched waveguides in TFLN next to each other are shown. We can see that the sidewalls of the waveguides are rough. This can be observed even better in the zoom in at the right top corner. This roughness of the sidewalls is due to the etching process and leads to an increase of losses. For the waveguide type analyzed in this thesis, there are two dominate types of losses. The first one are scattering losses and the second type is radiation or bending losses [34]. Absorption losses are always present, but in our case they are small and can be neglected [35]. Other kinds of losses that can occur in a waveguide are mode coupling losses as well as losses at interfaces

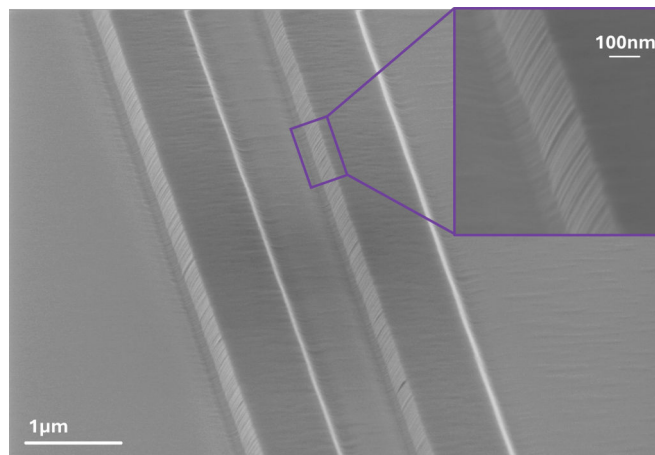


Fig. 2.2: The SEM image shows two etched waveguides in TFLN next to each other. We can see that the sides of the waveguides are rough due to the etching process. This is also shown in the close up on the right side. The roughness leads to increased propagation losses.

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e.g., light going from a straight waveguide into a bending, but these will not be discussed here.

Scattering Losses

Scattering losses can have different origins. The first one is scattering losses due to imperfections in the crystal structure. The second kind are due to roughness at the interfaces between film and substrate/cover. These type of losses can be understood in the ray optical picture. If the light bounces in the waveguide and there is any kind of discontinuity in the material, the light is reflected at this point which increases its angle and some part of the light is lost. This loss is a linear effect and can be described by the Lambert-Beer law [36]:

$$P = P_0 \cdot e^{-\alpha L} \quad (2.21)$$

$$\alpha = \frac{1}{L} \cdot \ln \left(\frac{P_0}{P} \right). \quad (2.22)$$

Here, P_0 denotes the input power of the light, α the attenuation coefficient in 1/cm and L the length, in this case here, of the waveguide. In this thesis, we will give the losses not in 1/cm but rather in dB/cm. The translation is as follows [34]:

$$10 \cdot \log_{10}(e) \cdot \alpha \left[\frac{1}{cm} \right] \approx 4.34 \cdot \alpha \left[\frac{1}{cm} \right] = \alpha \left[\frac{dB}{cm} \right]. \quad (2.23)$$

Radiation Losses

One way to look at radiation or bending losses is via the phase fronts of the propagating wave. For this, we imagine a wave inside a bending as depicted in Figure 2.3 (a). This wave has phase fronts perpendicular to the propagation direction and these need to be preserved while propagating in the bending. This, in turn, means that the outer part of the wave needs to travel faster than the inner part since the way is longer. At some point this increase in the velocity would result in a velocity faster than the speed of light. Since this is not possible, the light would be radiated and thus lost [34].

The other approach to radiation losses is the equivalent index profile [37, 38]. For this concept, we look again at the phase fronts. As already stated the phase fronts at the outer part need to travel faster. This can also be seen as a modified refractive index profile, since

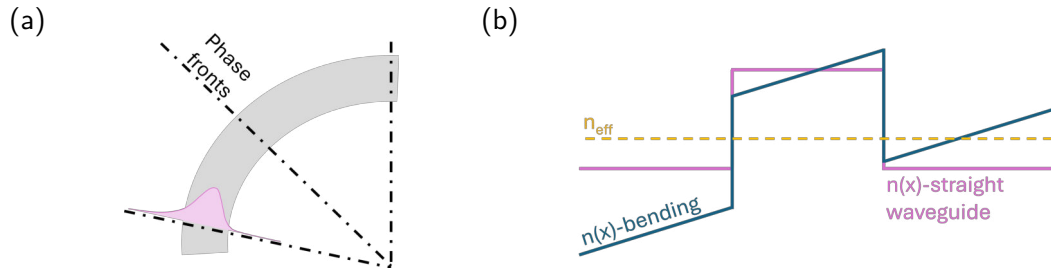


Fig. 2.3: Two approaches to look at radiation losses in waveguides. (a) Radiation losses can be explained via the phase fronts of the wave. (b) An equivalent index profile which is due the bending can also be used to describe bending losses.

an increase in velocity can also be seen as a decrease in the refractive index (see Figure 2.3 (b)). This means in the picture of the quantum mechanical finite potential well that the refractive index profile is tilted. A solution of this asymmetric index profile would be a guided mode inside the waveguide with an evanescent field. However at the higher tilt which corresponds to the outer part of the waveguide, the evanescent field can start to oscillate again. This oscillating evanescent field describes the radiation losses. [37, 38] The smaller the bending radius the more pronounced the tilt is. If the index profile has a large contrast, a stronger tilt which means a smaller bending is acceptable without bending losses. This means a waveguide with a high index contrast can have small bendings without high bending losses.

This ends the introduction into the basics of waveguides which are the foundation of integrated optics. Next, we will move on to more complex structures that are composed of waveguides such as directional couplers or resonators.

2.2 Integrated Devices

Integrated devices are composed of waveguides and can have various forms and functionalities. Here, we will look at two different devices that are crucial for integrated (quantum) optics circuits. The first one is a directional coupler which is a standard devices to combine or split light also used in modulator schemes. This device is also used to couple light to a resonator which we will discuss afterwards.

2.2.1 Directional Coupler

A directional coupler is a fundamental building block for many integrated devices such as filters, modulators and wavelength (de-) multiplexers [36]. A sketch of a directional coupler is shown in Figure 2.4. The name directional coupler means that the propagation direction of light is not changed in the coupler [34]. A directional coupler consists of two

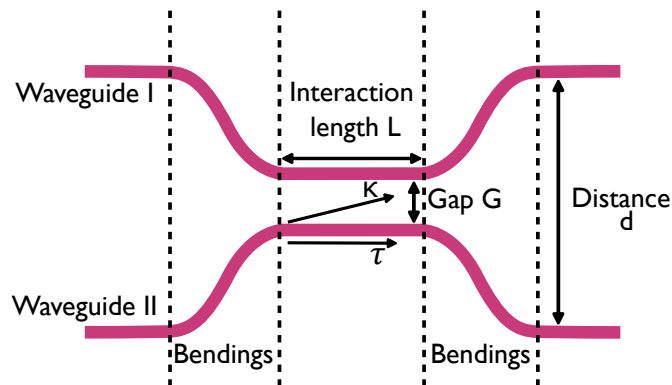


Fig. 2.4: Sketch of a directional coupler. Two waveguides are brought in close proximity to each other for an interaction length L separated by a gap G in which the light can couple between the waveguides. Afterwards the waveguides are separated again via bendings until they reach a distance d from each other. The coupling is described by the through and cross-port coefficients τ and κ , respectively.

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waveguides that are in close proximity to each other and optically coupled [36]. We call the upper waveguide waveguide I and the lower waveguide waveguide II. At the beginning of the coupler the two waveguides are well separated with a distance d and are then getting closer to each other via bendings until they are only separated by a small gap G . Then, they are parallel to each other for the interaction length L and separated again via bendings.

The coupling of light in such a structure can be described in the following way. The energy is transferred between the two waveguides due to the overlapping evanescent fields of the two waveguides if they are close enough to each other. Then, the evanescent field of one waveguide can reach into the second waveguide and can excite a mode there. [39] With increasing length of the coupler more and more energy is transferred from one waveguide to the other one [36]. Thus, the coupling is determined by the gap between the waveguides and the length of the coupling region [40].

To describe the behavior of such a directional coupler mathematically, the concept of supermodes or local normal modes is used [36, 39]. In this approach, the coupler is treated as a five layer system. For this system, one has to calculate the eigenmodes by solving the wave equation in each layer and enforcing continuity of the transverse electric and magnetic fields at the boundaries. Doing so leads to the supermodes of the layer system. In Figure 2.5 the first two supermodes are displayed. It can be seen that one mode has an even symmetry and the other one a odd one. Thus, the modes are called the symmetric and anti-symmetric mode. It is evident that the maximum of these modes are in the waveguide region. [39] From these two supermodes we can see that if we add both together, we only get a mode in the left waveguide. If we subtract the modes, we get only a mode in the right waveguide. Therefore, we can describe the modes in the following way

$$A_I = B_s + B_a \quad (2.24)$$

$$A_{II} = B_s - B_a, \quad (2.25)$$

where B_s and B_a are the amplitudes of the symmetric and anti-symmetric modes and A_I and A_{II} the amplitudes of the modes in the first and second waveguide [39]. If we

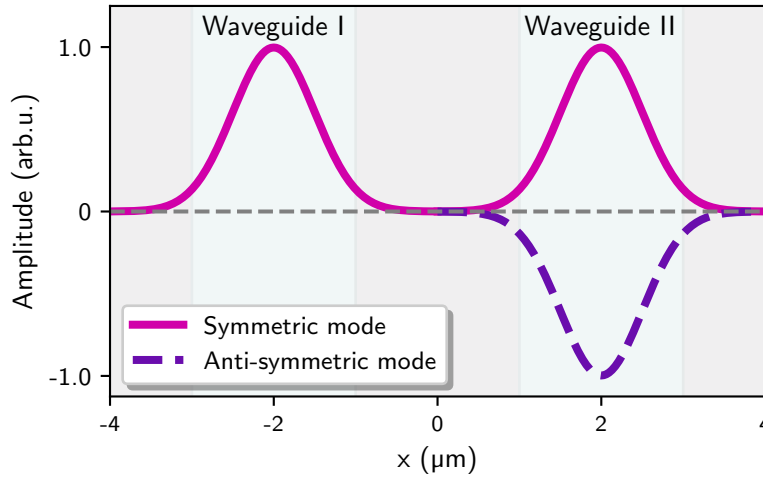


Fig. 2.5: Supermodes of a five layer system consisting of the two waveguides regions in light blue and the three regions next to the waveguides in grey. In such a system two modes can form with are called the symmetric and anti-symmetric mode.

2.2. Integrated Devices

neglect losses, we can presume that the symmetric and anti-symmetric mode travel along the waveguide with the following field distribution [39]:

$$B_s(z) = B_s(0) \cdot e^{i\beta_s z} \quad (2.26)$$

$$B_a(z) = B_a(0) \cdot e^{i\beta_a z}. \quad (2.27)$$

Here, β_s and β_a are the propagation constants of the symmetric and anti-symmetric mode. If we now couple light into our system, we would excite the symmetric and the anti-symmetric mode. Assuming we couple all the light in the first waveguide, we would have the following initial power distribution

$$A_I(0) = 1 \quad (2.28)$$

$$A_{II}(0) = 0, \quad (2.29)$$

which leads to [39]:

$$B_s(0) = B_a(0) = \frac{1}{2}. \quad (2.30)$$

Thus, all the light is in the first waveguide and both supermodes are excited. While the two supermodes travel along the waveguide, they collect a phase shift, since the two propagation constants are different [36]:

$$\varphi = \int_0^L (\beta_s(z) - \beta_a(z)) \, dz. \quad (2.31)$$

Assuming the propagation constants are independent of z , the integral yields [36]:

$$\varphi = (\beta_s - \beta_a) \cdot L. \quad (2.32)$$

If this phase shift is equal to π , all the light will couple from one waveguide to the other one. This is also called the crossover length L_x [40].

The power in each waveguide can also be calculated by evaluating the modes in each arm together with the phase shift [36, 40]:

$$P_I \sim A_I A_I^* \rightarrow \tau^2 = \frac{P_I}{P_0} = \cos^2 \left(\frac{\Delta\beta}{2} L \right) \quad (2.33)$$

$$P_{II} \sim A_{II} A_{II}^* \rightarrow \kappa^2 = \frac{P_{II}}{P_0} = \sin^2 \left(\frac{\Delta\beta}{2} L \right). \quad (2.34)$$

Here, τ is the through-port coefficient and κ the cross-port coefficient [40]. The coupling coefficient C as well as the crossover length L_x can also be expressed in terms of the refractive index via the relation $\beta = \frac{2\pi n}{\lambda}$ [40]:

$$C = \frac{\Delta\beta}{2} = \frac{\pi \Delta n}{\lambda} \quad (2.35)$$

$$L_x = \frac{\lambda}{2\Delta n}. \quad (2.36)$$

Here, $\Delta\beta = \beta_s - \beta_a$ is the difference of the propagation constants of the symmetric and anti-symmetric mode. Consequently, Δn is the difference of the effective refractive index of the symmetric and anti-symmetric mode. This means, the coupling depends on the effective refractive index of the two modes which are influenced by the gap as well as the waveguide geometry. In Figure 2.6 the coupling behavior is shown. We can see that the light oscillates between the two waveguides. This coupling between two waveguides can also be utilized to couple light to various waveguide structures such as a waveguide forming a loop, creating a resonator. Thus, next we discuss resonators.

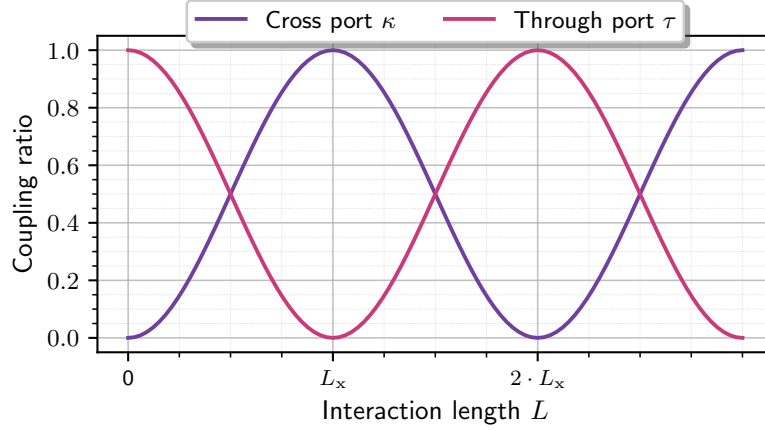


Fig. 2.6: Coupling ratio as a function of the length of the coupling region. The light oscillates between the two waveguides and every crossover length L_x the light completely coupled from one waveguide to the other.

2.2.2 Resonators

Optical resonators are crucial components in the field of integrated optics since they can be utilized for many different applications such as filters, laser cavities, etc.. A resonator confines light of a specific frequency which is determined by the structure of the resonator [32]. There are many different ways of constructing an optical resonator. Just two parallel partially reflecting mirrors are already an optical resonator which is called a Fabry-Pérot resonator. However, there are also more complex resonator structures such as whispering gallery resonator or photonic crystals [32, 39]. Also, when we make a loop out of one waveguide, we create an optical resonator called ring resonator. To couple light to the resonator, we can utilize a directional coupler. The basic functionality of all these different types can be determined from the most easiest resonator which is the Fabry-Pérot resonator. Thus, in this chapter we will first discuss the principle behind this resonator type and move then on to an integrated ring or racetrack resonator.

Fabry-Pérot Resonator

A Fabry-Pérot resonator consists of two plain, parallel and partly reflective mirrors with the reflection coefficient r and the transmission coefficient t which are at the distance d from each other [41]. To understand how this resonator works, we assume that we have one initial wave E_0 that hits the first mirror. A part of the wave is reflected from the mirror and a part is transmitted. The transmitted part is now between the two mirrors. Here, the wave accumulates a phase of $e^{i\varphi}$ with $\varphi = 2 \cdot d \cdot \beta$ and experiences round trip losses $e^{-\alpha d}$ with α as the loss coefficient. When it hits the second mirror again a part of the light is reflected and a part is transmitted. Thus, the initial beam is reflected back and forth between the two mirrors such that there are multiple waves inside the resonator [42]. The wave inside the resonator can now interfere with each other. If they interfere constructively, the intensity is enhanced. However, if they interfere destructively, the intensity is suppressed. The waves interfere constructively, if the optical path which is given by $2 \cdot n \cdot d$ with n the refractive index of the material between the two mirrors is a multiple of the wavelength λ of the wave

[33]:

$$2 \cdot n \cdot d = k \cdot \lambda. \quad (2.37)$$

From this condition, we can already conclude that only for specific wavelengths the output spectrum of a Farby-Pérot resonator will have resonances and for all other wavelengths we will witness destructive interference. The spectrum of a Fabry-Pérot resonator can be derived by summing up all the transmitted partial waves which yields using the harmonic series the following equation [36]:

$$I_t = \left\| \frac{E_t}{E_0} \right\|^2 = \frac{t_1^1 \cdot t_2^2 \cdot e^{-\alpha d}}{1 + r_1^2 \cdot r_2^2 \cdot e^{-\alpha d} - 2 \cdot r_1 \cdot r_2 \cdot e^{-\alpha d} \cdot \cos(\varphi)}. \quad (2.38)$$

A similar calculation for the reflected fields yields [36]:

$$I_r = \left\| \frac{E_r}{E_0} \right\|^2 = \frac{r_1^2 + r_2^2 \cdot e^{-2\alpha d} - 2 \cdot r_1 \cdot r_2 \cdot e^{-\alpha d} \cdot \cos(\varphi)}{1 + r_1^2 \cdot r_2^2 \cdot e^{-\alpha d} - 2 \cdot r_1 \cdot r_2 \cdot e^{-\alpha d} \cdot \cos(\varphi)}. \quad (2.39)$$

In Figure 2.7 the reflected and transmitted powers are shown. As already discussed in the beginning, we have distinct resonances at multiple of $2 \cdot \pi$ due to the accumulated phase inside the resonator. If the reflected intensity shows a minimum, the transmitted spectrum shows a maximum. If the resonance condition is not fulfilled, the transmitted intensity drops and the reflected one increases.

From this spectrum, we can now determine two important parameters of a resonator. The first parameter is the free spectral range (FSR). This parameter describes the distance between two adjacent peaks which is given by [41]:

$$\text{FSR} = \frac{\lambda^2}{2 \cdot n_g \cdot d}. \quad (2.40)$$

Here, n_g is the group index. The FSR is inversely proportional to the optical path length between the mirrors. Thus, the larger the path length, the smaller the FSR.

The second parameter we want to discuss is the resonator's finesse. The finesse is specified

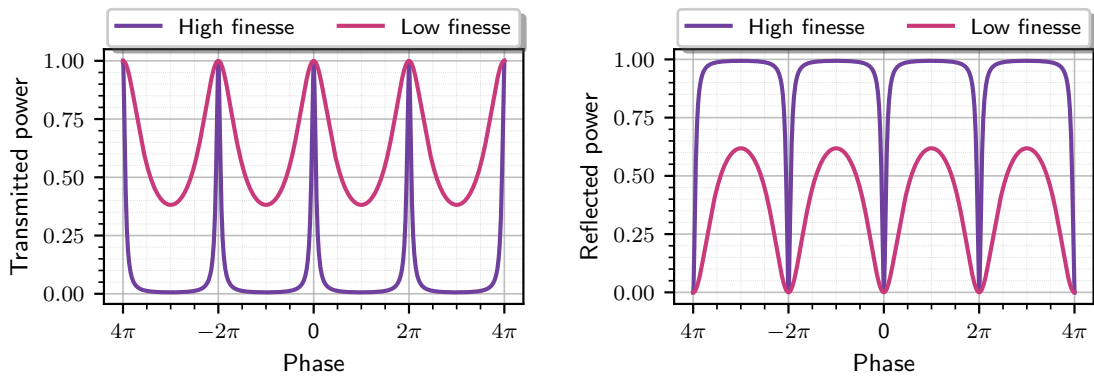


Fig. 2.7: (a) Transmitted and (b) reflected spectrum of a lossless ($\alpha=0$) Fabry-Pérot resonator with a high and low finesse.

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as the ratio between the FSR and the full width at half maximum (FWHM) of the resonance peaks. This FWHM is given by [43]:

$$\text{FWHM} = \frac{(1 - r_1 \cdot r_2 \cdot e^{-\alpha d/2})}{r_1^2 \cdot r_2^2 \cdot e^{-\alpha d} \cdot \pi \cdot n_g \cdot d}. \quad (2.41)$$

From this follows that the finesse is:

$$\text{Finesse} = \frac{\text{FSR}}{\text{FWHM}} = \frac{\pi \cdot r_1^2 \cdot r_2^2 \cdot e^{-\alpha d}}{1 - r_1 \cdot r_2 \cdot e^{-\alpha d/2}}. \quad (2.42)$$

If the finesse is high, the spectrum shows sharp peaks (see Figure 2.7 purple lines) [32]. For a high finesse, the mirrors need to be highly reflective and the losses need to be low. If the resonator has low reflectivities and/or high losses the finesse is low and thus, the resonance peaks broaden (see Figure 2.7 magenta lines). The finesse intuitively describes the number of round trips the light experiences inside the resonator. This kind of resonator can be used to estimate the propagation losses of waveguides which we will discuss next.

Estimation of Propagation Losses using a Fabry-Pérot Resonator

The Fabry-Pérot method [44] is based on the fact that a waveguide sample with its end faces can be seen as a Fabry-Pérot resonator. In this case, the end faces correspond to the mirrors of the resonator. By measuring the contrast of the transmission spectrum of the resonator, the propagation losses can be estimated, see Figure 2.8. To derive the equation for the loss calculation, we start with the transmitted intensity, see Equation 2.38. We multiply the equation by the coupling efficiency η since not all the input light I_0 is coupled to the waveguide. This yields:

$$I_T = \frac{t_1^2 t_2^2 e^{-\alpha L}}{1 + r_1^2 r_2^2 e^{-\alpha L} - 2r_1^2 r_2^2 e^{-\alpha L} \cos(\varphi)} I_0 \eta. \quad (2.43)$$

Moreover, we assume that the transmittivity as well as the reflectivity of both end faces are the same ($R = r_1^2 = r_2^2, T = t_1^2 = t_2^2$) which leads to:

$$I_T = \frac{T^2 e^{-\alpha L}}{1 + R^2 e^{-\alpha L} - 2R e^{-\alpha L} \cos(\varphi)} I_0 \eta. \quad (2.44)$$

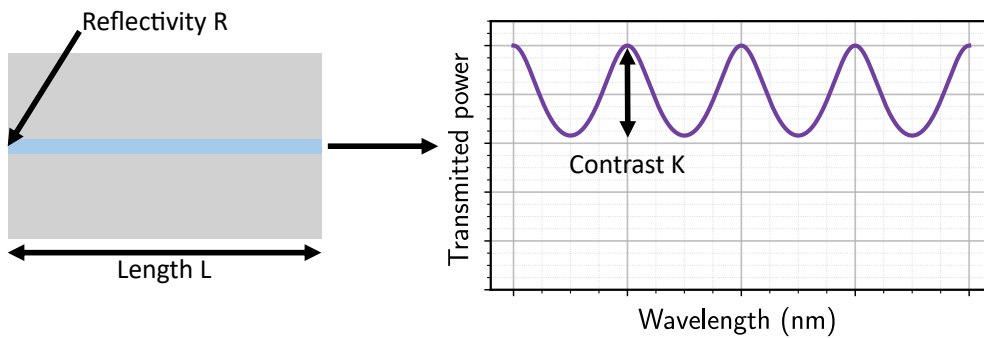


Fig. 2.8: Conceptual drawing of the Fabry-Pérot loss measurement technique.

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With the identity $\cos(2x) = 1 - 2\sin^2(x)$ and by introducing the combined loss-reflection factor $\tilde{R} = Re^{-\alpha L}$, we get:

$$I_T = \frac{T^2 e^{-\alpha L}}{(1 + \tilde{R})^2 + 4\tilde{R} \sin^2(\varphi/2)} I_0 \eta. \quad (2.45)$$

This shows that the transmitted intensity varies periodically when changing the phase which is determined by the optical path length. Thus, by for example heating the sample or changing the wavelength the phase can be altered. Then, the losses can be estimated from the contrast K which is the ratio of the difference of the maximal and minimal intensity and the sum of both:

$$K = \frac{I_{\max} - I_{\min}}{I_{\max} + I_{\min}}. \quad (2.46)$$

Since $I_{\max}$ is given when $\sin^2(\varphi/2) = 0$ and $I_{\min}$ when $\sin^2(\varphi/2) = 1$, the contrast K can also be written as:

$$K = \frac{2\tilde{R}}{1 + \tilde{R}^2} \quad (2.47)$$

which is independent of the coupling efficiency. Hence, $\tilde{R}$ is given by:

$$\tilde{R} = \frac{1}{K} (1 - \sqrt{1 - K^2}). \quad (2.48)$$

If we now assume that K is small which is given for low-finesse resonators, $\sqrt{1 - K^2}$ can be approximated via Taylor series $1 - K^2/2$ and thus $\tilde{R} \approx K/2$. If we use $\tilde{R} = Re^{-\alpha L}$ and solve for α to estimate the losses we arrive at [44]:

$$\alpha \approx \frac{4.34}{L} (\ln(R) + \ln(2) - \ln(K)). \quad (2.49)$$

The pre-factor 4.34 ($= 10 \cdot \log_{10}(e)$) is due to the fact that the losses are given in dB/cm. With the equation the propagation losses of single straight waveguides can be estimated. For this, the transmission spectrum is measured which is used to determine the contrast K . Moreover, we need to know the length L of the sample and the reflectivity R of the end faces of the waveguides precisely. In the weak light confinement conditions, the end face reflectivity is calculated using Fresnel law [45]:

$$R = \left(\frac{1 - n}{1 + n} \right)^2. \quad (2.50)$$

Here, n is the refractive index of the material as for example in our case lithium niobate. A precise knowledge of the end face reflectivity of the investigated waveguides is crucial for a precise loss estimation. This method has the benefit of being able to measure the losses of single straight waveguides which are the simplest building block for integrated circuits. Also this method can be facilitated in ready-to-use devices such as phase modulators, single pass photon sources or frequency mixers.

Now, we can understand the principle of a Fabry-Pérot resonator and know the most important parameters which characterize a resonator's spectrum. Moreover, we also understand how such a resonator can be used to measure the propagation losses of waveguides. Thus, we can use our gained knowledge to discuss ring or racetrack resonators.

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Racetrack Resonator

A racetrack resonator is a special form of a ring resonator. The layout of it is shown in Figure 2.9. It consists of a waveguide connected in a loop, with two bendings and two straight section, and a bus waveguide that routes the light towards the racetrack. In the coupling region the bus waveguide is only separated by a small gap G from the racetrack. In

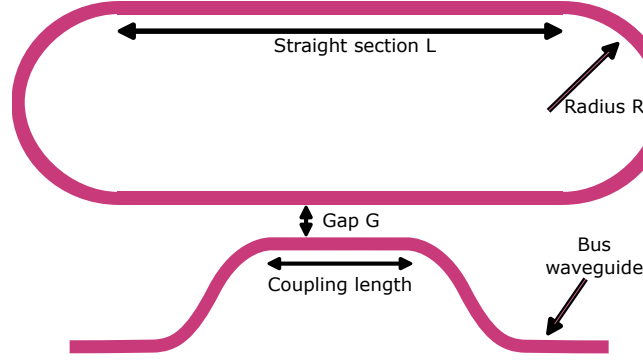


Fig. 2.9: Sketch of a racetrack resonator connected to a bus waveguide. A racetrack resonator consists of two straight parts and two bendings. The light couples from the bus waveguide to the racetrack resonator and back.

this area light can couple between the bus waveguide and the racetrack in the same fashion as a directional coupler.

To mathematically describe the output spectrum, we can go back to a Fabry-Pérot resonator. We can use Equation 2.39 and assume that the second mirror has a 100 % reflectivity. Hence, the output intensity can be described via the following equation [43]:

$$I_{RR} = \frac{r^2 + e^{-\alpha L} - 2re^{-\alpha L/2}\cos(\varphi)}{1 - 2e^{-\alpha L/2}r\cos(\varphi) + re^{-\alpha L}}. \quad (2.51)$$

Therefore, the spectrum looks similar to a Fabry-Pérot resonator and we get again resonances when the optical length of the racetrack is equal to multiple of the wavelength. The depth of the resonance peaks depends on the coupling coefficient of the directional coupler and the losses of the racetrack resonator. Here, we differentiate between 3 cases which are critical coupling, overcoupling and undercoupling.

Critical coupling is present if the through port coefficient τ is the same as the round trip losses ($\tau = e^{-\alpha L/2}$). This means that the amplitude of the light that traveled once around the racetrack is the same as the transmitted light in the bus waveguide. So, the external coupling rate matches the intrinsic loss rate. Thus, when the light couples back from the ring to the bus waveguide, it experiences a π phase shift. Now, we have two fields with opposite signs but the same amplitude. This leads to a complete destructive interference and no light is transmitted.

For overcoupling the through port coefficient τ is smaller than the losses ($\tau < e^{-\alpha L/2}$) and for undercoupling larger than the losses ($\tau > e^{-\alpha L/2}$). This leads to a field in the racetrack with a higher or lower amplitude which leads to a non-zero transmitted intensity. [43]

Figure 2.10 shows the three spectra for critical, over- and undercoupling. The condition of critical coupling leads to no transmitted light at resonance whereas over- and undercoupling have a non-zero transmitted intensity at resonance. Otherwise, a racetrack resonator can be

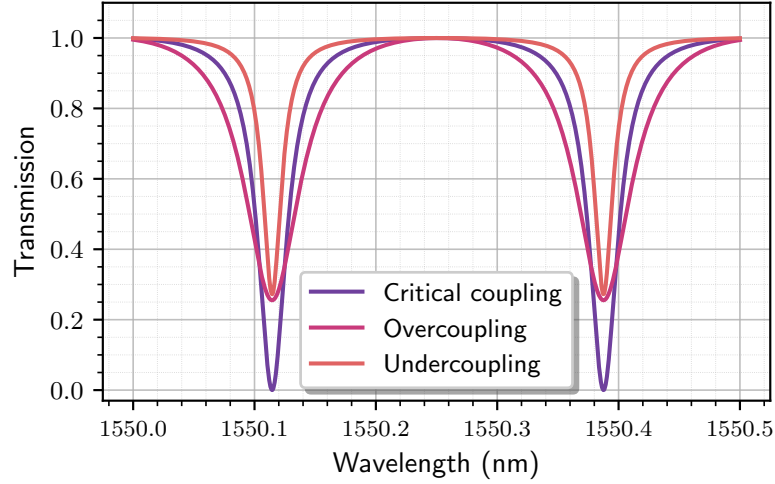


Fig. 2.10: Output spectrum of a racetrack resonator for the three different coupling cases which depend on the relation between the losses of the resonator and the coupling to the resonator. If the resonator is critically coupled there is no transmission at resonance. If it is under- or overcoupled, light is transmitted at resonance.

characterized with the same parameters as a Fabry-Pérot resonator. The FSR of a racetrack resonator is given by [43]:

$$\text{FSR} = \frac{\lambda^2}{n_g \cdot L}. \quad (2.52)$$

Additionally, the FWHM is defined by [43]:

$$\text{FWHM} = \frac{(1 - ra)\lambda^2}{\pi n_g L \sqrt{re^{-\alpha L/2}}}. \quad (2.53)$$

Another very useful parameter that we not yet introduced is the quality factor or short Q-factor. The Q-factor is determined by the relation between the energy inside the racketrack and the losses in each round trip [40]:

$$Q = \omega_0 \frac{\text{Energy stored}}{\text{Power loss per cycle}}. \quad (2.54)$$

There are different kinds of Q-factors. The first one is the loaded Q-factor Q_L which takes the coupling to the bus waveguide into account. The second kind is the unloaded Q-factor Q_U which does not consider the coupling to the bus waveguide and is solely determined by the losses of the resonator [43].

Another definition of the Q-factor is given by the ration between the wavelength of the resonance peak and the sharpness of the reson the resonance peak, e.g., the FWHM [46]:

$$Q_L = \frac{\lambda_{\text{res}}}{\text{FWHM}} = \frac{\pi n_g L \sqrt{re^{-\alpha L/2}}}{(1 - re^{-\alpha L/2})\lambda}. \quad (2.55)$$

This equation describes the loaded Q-factor. Accordingly, we can calculate the unloaded Q-factor by assuming that the reflectivity is 100 % ($r = 1$) [46]:

$$Q_U = \frac{\pi n_g L e^{-\alpha L/4}}{(1 - e^{-\alpha L/2})\lambda}. \quad (2.56)$$

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We can simplify this further by using a Taylor expansion of $e^{-\alpha L/2}$ [46]:

$$e^{-\alpha L/2} = 1 + (-\alpha L/2) = 1 - \alpha L/2. \quad (2.57)$$

If we insert this into Equation 2.56 we get:

$$Q_U = \frac{\pi n_g L e^{-\alpha L/4}}{(1 - 1 + \alpha L/2)\lambda} = \frac{\pi n_g (1 - \alpha L/2)}{\alpha L/2\lambda} = \frac{2\pi n_g}{\alpha \lambda_0} (1 - \alpha L/4). \quad (2.58)$$

Furthermore, we assume that $\alpha L \ll 1$ for a low loss racetrack. This leaves us with [46]:

$$Q_U = \frac{2\pi n_g}{\alpha \lambda_0}. \quad (2.59)$$

Now, we have an expression for the loaded and the unloaded Q-factor. In the experiment, we measure the loaded Q-factor by determining the FSR and FWHM of the resonance peaks, since we measure the complete racetrack response including the bus waveguide. However, the unloaded Q-factor is the more interesting figure of merit since it only depends on the losses of the resonator and we can therefore estimate the propagation losses via the unloaded Q-factor. Thus, we want a relation between the loaded and the unloaded Q-factor with which we can directly estimate the unloaded Q-factor from our measured racetrack spectrum.

To derive such a relation, we start with the definition of the Q-factor as shown in Equation 2.54. Since the Q-factor is determined by the energy inside the racetrack in relation to the losses the light experiences each round trip, it is crucial to understand the different loss mechanism inside the racetrack. On the one hand, we have intrinsic losses of the waveguide due to material absorption, scattering losses etc.. This we will describe by the decay rate $1/\tau_0$ due to intrinsic losses. On the other hand, energy from the ring is lost due to the coupling to the bus waveguide. This loss mechanism is described by the decay rate $1/\tau_c$ due to coupling. Since the loaded Q-factor takes intrinsic losses as well as coupling losses into account, we can rewrite the loaded Q-factor as follows:

$$Q_L = \frac{\omega_0}{(1/\tau_0 + 1/\tau_c)}. \quad (2.60)$$

As mentioned above, the unloaded Q-factor describes the racetrack without coupling and can thus be defined as:

$$Q_U = \frac{\omega_0}{1/\tau_0}. \quad (2.61)$$

Consequently, we can find a relation between the loaded and unloaded Q-factor [40]:

$$\frac{1}{Q_L} = \frac{1}{Q_U} + \frac{1}{Q_C}. \quad (2.62)$$

Here, Q_C is the coupling Q-factor which is solely determined by the coupling losses. With this relation we can still not determine the unloaded Q-factor from a measured racetrack spectrum. Therefore, we derive another important parameter which is the reflection coefficient. Since a racetrack resonator behaves like a lossy, external driven harmonic oscillator, e.g. an RCL circuit, we can relate to the general harmonic oscillator theory to derive the reflection coefficient which is then given by [47, 48]:

$$R(\omega_0) = \left| \frac{1 - \frac{\tau_c}{\tau_0}}{1 + \frac{\tau_c}{\tau_0}} \right|^2. \quad (2.63)$$

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This parameter describes the peak depth at resonance and is thus easily measured. If we solve this for $\frac{\tau_c}{\tau_0}$ we get [48, 46]:

$$\frac{\tau_c}{\tau_0} = \frac{1 - \sqrt{R(\omega_0)}}{1 + \sqrt{R(\omega_0)}}. \quad (2.64)$$

Now, we go back to Equation 2.62 and rewrite it:

$$Q_L = \frac{Q_U Q_C}{Q_C + Q_U} = \frac{Q_U}{1 + \frac{Q_U}{Q_C}} = \frac{Q_U}{1 + \frac{\tau_c}{\tau_0}}. \quad (2.65)$$

In the last step, we use the relation between the Q-factors and the decay rates. Now we can insert Equation 2.64 into this equation [46, 49]:

$$Q_L = Q_U \cdot \left(1 + \frac{1 - \sqrt{R(\omega_0)}}{1 + \sqrt{R(\omega_0)}} \right)^{-1} \quad (2.66)$$

$$Q_U = \frac{2 \cdot Q_L}{1 \pm \sqrt{R(\omega_0)}}. \quad (2.67)$$

The sign depends on the coupling condition of the resonator. If the resonator is undercoupled, the plus sign has to be used. If it is overcoupled, the minus sign. With this equation, it is now possible to estimate the unloaded Q-factor when measuring the loaded Q-factor. From the unloaded Q-factor it is then possible to determine the propagation losses of the racetrack. The exact measurement method to do this will be discussed next.

Estimation of Propagation Losses using Racetrack Resonators

The racetrack resonator method is based on the relation of the Q-factor and the losses of a racetrack resonator. By estimating the Q-factor of multiple racetrack resonators with different lengths of the straight parts, a correlation between the propagation losses and racetrack length can be found [18]. To estimate the losses with the racetrack method, we need at least three racetrack resonators with the same bending radius. Only the length of the straight part changes. Keeping the bending the same insures that an increase in losses is only due to the increase in the length of the resonator. Thus, by estimating the Q-factor of at least three different racetrack resonators, the propagation losses can be derived by the increase of the total losses with increasing racetrack length.

To derive the equation for calculating the total losses of a racetrack resonator we start with Equation 2.59 and solve it for α :

$$\alpha = \frac{2\pi n_g}{\lambda_0 Q_u}. \quad (2.68)$$

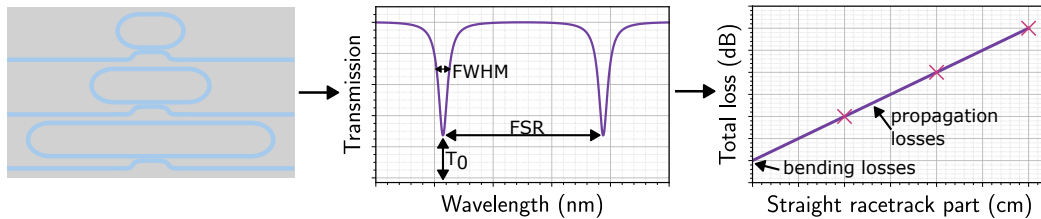


Fig. 2.11: Conceptual drawing of the racetrack resonator method to measure the propagation losses.

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Next, we use Equation 2.52 to substitute the group index n_g and Equation 2.67 to replace the unloaded Q-factor:

$$\alpha = \frac{\pi \lambda_0}{\text{FSR } L Q_L} \left(1 \pm \sqrt{R(\lambda_0)} \right). \quad (2.69)$$

Lastly, we use Equation 2.55 as the definition of the Q-factor which yields the following equation [50]:

$$\alpha = \frac{4.34 \pi \text{ FWHM}}{L \text{ FSR}} \left(1 \pm \sqrt{R(\lambda_0)} \right). \quad (2.70)$$

Again, the factor 4.34 is due to the unit dB/cm and $\pm$ depends if the resonator is under- or overcoupled, as discussed before. Thus, to calculate the total losses of a racetrack resonator, we need to measure the full width at half maximum FWHM, the free spectral range FSR, the length of the sample L, the resonance wavelength λ_0 as well as the transmitted power at resonance $R(\lambda_0)$ from the measured transmission spectrum. To decouple the propagation losses from the bending losses, we first multiply the losses by the total length of the resonator to get the total losses in dB. Next, we plot the total losses as a function of the straight length of the resonator, as shown in Figure 2.11. From a linear fit, we extract the propagation losses as the slope and the bending losses as the y-intersection. However, when the bending losses are negligible also a single racetrack or ring resonator can be used to estimate the losses.

The advantage of this method is that in contrast to the Fabry-Pérot method no external parameter is needed. However, the disadvantages are that with this method we cannot measure the losses of single straight waveguides which is for example important when characterizing single pass single photon sources. Moreover, we need to measure multiple racetracks for which we need to ensure a fabrication homogeneity over the sample.

This concludes the section about waveguides and integrated devices which are the key components of every integrated optical circuit. However, an integrated (quantum) optical circuit does not solely consists of waveguides, linear devices and effects. Nonlinear optical effects are often utilized to generate light at a different frequency or modulate the light on chip. Thus, next we will discuss the principle of nonlinear optics.

2.3 Nonlinear Optics

Up to now, we discussed waveguide structures based on linear optics. However, this thesis deals with integrated quantum optics in TFLN waveguides. Next to waveguides structures, also nonlinear effects are utilized to generate light at different frequencies. Thus, in this chapter, we discuss the basic principles of nonlinear optics. Here, we first focus on the nonlinear material response to understand the underlying principle. Next, we derive the nonlinear wave equation and introduce the sum frequency generation process as well as phase matching.

2.3.1 Nonlinear Material Response

In the field of nonlinear optics we experience effects in which the optical properties of the material can be altered by applying an electric field [51]. To describe the effects we use

2.3. Nonlinear Optics

the polarization $P(t)$ which is defined as the dipole moment per unit volume [51]. In linear optics, we limit ourselves only to linear effects, where the electric field does not influence the material properties and its response - the polarization - is linear with the electric field [52]:

$$P(t) = \epsilon_0 \cdot \chi^{(1)} E(t). \quad (2.71)$$

Here, ϵ_0 is the electric field constant and $\chi^{(1)}$ is the linear susceptibility. In the case of nonlinear optics the polarization does not only depend linear on the electric field. If the applied electric field is strong enough, also higher order terms can occur as such a strong field causes perturbations in the material. Thus, we express the polarization as a power series in the field strength $E(t)$ [51]:

$$P(t) = \epsilon_0 \left(\chi^{(1)} E(t) + \chi^{(2)} E(t)^2 + \chi^{(3)} E(t)^3 + \dots \right). \quad (2.72)$$

Here, $\chi^{(2)}$ and $\chi^{(3)}$ are the second and third order susceptibility. In this thesis, we focus on second order nonlinearities and the processes that originate from this, which can be described via the following term:

$$P^{(2)}(t) = \epsilon_0 \chi^{(2)} E(t)^2. \quad (2.73)$$

To analyze which processes can arise due to a second order susceptibility, we assume two waves E_1 and E_2 with frequencies ω_1 and ω_2 . For this, the input fields are monochromatic waves and can be described as follows [51]:

$$E(t) = E_1 e^{-i\omega_1 t} + E_2 e^{-i\omega_2 t} + cc. . \quad (2.74)$$

Inserting this into Equation 2.73 yields:

$$\begin{aligned} P^{(2)}(t) = \epsilon_0 \chi^{(2)} & (E_1^2 e^{-i2\omega_1 t} + E_2^2 e^{-i2\omega_2 t} + 2E_1 E_2 e^{-i(\omega_1 + \omega_2)t} \\ & + 2E_1 E_2 e^{-i(\omega_1 - \omega_2)t} + E_1 E_1^* + E_2 E_2^*). \end{aligned} \quad (2.75)$$

We can identify four different processes that oscillate with different frequencies and two static processes. For better visibility, we write $P^{(2)}(t)$ as a sum of its Fourier components [51]:

$$P^{(2)}(t) = \sum_n P^{(2)}(\omega_n) e^{-i\omega_n t}. \quad (2.76)$$

With this we can identify the following frequency components and names of the processes [51]:

$$P(2\omega_1) = \epsilon_0 \chi^{(2)} E_1^2 \text{ second harmonic generation of } \omega_1 \quad (2.77)$$

$$P(2\omega_2) = \epsilon_0 \chi^{(2)} E_2^2 \text{ second harmonic generation of } \omega_2 \quad (2.78)$$

$$P(\omega_1 + \omega_2) = 2\epsilon_0 \chi^{(2)} E_1 E_2 \text{ sum frequency generation} \quad (2.79)$$

$$P(\omega_1 - \omega_2) = 2\epsilon_0 \chi^{(2)} E_1 E_2 \text{ difference frequency generation} \quad (2.80)$$

$$P(0) = 2\epsilon_0 \chi^{(2)} (E_1 E_1^* + E_2 E_2^*) \text{ optical rectification.} \quad (2.81)$$

The first two processes are called second harmonic generation (SHG) for which the frequency of the generated field is twice the input frequency. Sum frequency generation (SFG) describes

Chapter 2. Foundations and Theoretical Concepts

a process in which the generated field oscillates with the sum of the two input frequencies. In contrast, for a difference frequency generation (DFG) the difference of the two input frequencies results in the output frequency. The last process that can occur is optical rectification (OR) in which a static electric field is generated. These are the four different processes that can take place due to a material with a second order susceptibility. However, not all materials exhibit a second order susceptibility. Only non-centro symmetric crystals have a $\chi^{(2)}$ and are called second order nonlinear materials. To understand this, we discuss the case of a centro symmetric crystal. If we apply an electric field to a centro symmetric crystal which introduces a nonlinear polarization, this polarization needs to be the same if we rotate the crystal, since we assumed our crystal has inversion symmetry. If the field components change the sign also the nonlinear polarization has to switch the sign. Looking at Equation 2.73 it can be seen that if the fields change the sign, the polarization does not change it. Thus, this condition is only fulfilled if all elements of the second order susceptibility are zero. [51, 53] Consequently, a second order nonlinearity only appears in non-centro symmetric crystals. To describe the response of such a second order nonlinear material to applied electric fields in detail, we need to take into account that the nonlinear response of the material may depend on the direction of the applied field in relation to the crystal axis. Moreover, a nonlinear polarization can be induced in a different direction than the applied field. This means, we have to generalize Equation 2.73 to allow of an anisotropic response of the material. [53] Therefore, we can write the second order polarization in the following way [52]:

$$P_i^{(2)} = \epsilon_0 \cdot \sum_{j,k} \chi_{i,j,k}^{(2)} \cdot E_j \cdot E_k. \quad (2.82)$$

Now, the second order susceptibility is a second order tensor and i, j, k correspond to the Cartesian axes x, y, z of the crystal. This means that the tensor has 27 different components. However, due to the high symmetry of crystals many quantities are zero. [53] First of all, some of the 27 components are the same due to the reason that the material response cannot depend on the mathematical order of the fields for example $\chi_{x,y,z}^{(2)} E_y E_z = \chi_{x,z,y}^{(2)} E_z E_y$. This leaves us with only 18 independent components. [53] This is also known as the Kleinman symmetry condition which states that if the frequency of the involved optical fields are much smaller than the resonance frequency of the material, the nonlinear susceptibility is independent of the frequency. This leads to the fact that we can permute the indices of the susceptibility without having to permute the frequencies. [51]

Due to the Kleinman symmetry condition we can introduce a contracted notation. For this we use the following relation [51]:

$$d_{i,j,k} = \frac{1}{2} \chi_{i,j,k}^{(2)}. \quad (2.83)$$

Now, due to the Kleinman symmetry condition we can interchange the last two indices. Thus, we can introduce a contracted notation by adding j, k to l with the following prescription [51]:

$$l = \begin{cases} 1 & \text{for } jk=11 \\ 2 & \text{for } jk=22 \\ 3 & \text{for } jk=33 \\ 4 & \text{for } jk=23 \text{ or } 32 \\ 5 & \text{for } jk=13 \text{ or } 31 \\ 6 & \text{for } jk=12 \text{ or } 21. \end{cases} \quad (2.84)$$

2.3. Nonlinear Optics

With this the nonlinear susceptibility tensor can be written as the following 3x6 matrix [51]:

$$d_{i,l} = \begin{bmatrix} d_{11} & d_{12} & d_{13} & d_{14} & d_{15} & d_{16} \\ d_{21} & d_{22} & d_{23} & d_{24} & d_{25} & d_{26} \\ d_{31} & d_{32} & d_{33} & d_{34} & d_{35} & d_{36} \end{bmatrix}. \quad (2.85)$$

If we now use again the Kleinman symmetry condition and permute the indices freely, we can see that a few tensor elements are the same as for example $d_{16} = d_{121} = d_{211} = d_{21}$. Therefore instead of 18 independent elements we only have 10. Thus, we can now describe the second order nonlinear polarization in the following way [51]:

$$\begin{bmatrix} P_x(\omega_3) \\ P_y(\omega_3) \\ P_z(\omega_3) \end{bmatrix} = M \cdot \epsilon_0 \cdot \begin{bmatrix} d_{11} & d_{12} & d_{13} & d_{14} & d_{15} & d_{16} \\ d_{16} & d_{22} & d_{23} & d_{24} & d_{14} & d_{12} \\ d_{15} & d_{24} & d_{33} & d_{23} & d_{13} & d_{14} \\ E_x(\omega_1)E_x(\omega_2) \\ E_y(\omega_1)E_y(\omega_2) \\ E_z(\omega_1)E_z(\omega_2) \\ E_y(\omega_1)E_z(\omega_2) + E_z(\omega_1)E_y(\omega_2) \\ E_z(\omega_1)E_x(\omega_2) + E_x(\omega_1)E_z(\omega_2) \\ E_x(\omega_1)E_y(\omega_2) + E_y(\omega_1)E_x(\omega_2) \end{bmatrix}. \quad (2.86)$$

We can differentiate between three different types of processes depending on the polarization of the involved fields. If all three fields have the same polarization we will call it type 0. If the polarization of the generated field is different from the two applied fields we call it type 1. Additionally, type 2 describes if the two applied electric fields have different polarization. Here, M is a degeneracy factor and depends on the number of distinct permutations of the frequencies of the applied fields [51]. M is equal to 1 for type 0/1 SHG processes and M is equal to 2 for all other processes.

Until now we haven't seen that an electric field can induce a nonlinear material polarization. Next, we will investigate how the three interacting fields behave while propagating through the material.

2.3.2 Nonlinear Wave Equation

To study the propagation behavior of a coupling between the electric field and the induced polarization, we first derive the wave equation that describes the propagation of light through a nonlinear medium.

Thus, we start again with Maxwell's equations for a non-magnetic and non-conducting material:

$$\nabla \times \vec{E} = -\frac{\partial \vec{B}}{\partial t} \quad (2.87)$$

$$\nabla \times \vec{B} = \mu_0 \frac{\partial \vec{D}}{\partial t} \quad (2.88)$$

$$\nabla \cdot \vec{D} = 0 \quad (2.89)$$

$$\nabla \cdot \vec{B} = 0. \quad (2.90)$$

In contrast to the linear case which we discussed in the beginning, we assume now a nonlinear dependence of the polarization on the electric field to incorporate possible effects which

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arise due to a nonlinear polarization [51]:

$$\vec{D} = \epsilon_0 \vec{E} + \vec{P}. \quad (2.91)$$

To get to a nonlinear wave equation, we will apply a curl operator to Equation 2.87 similar as we did it for the linear wave equation. Doing so yields [51]:

$$-\nabla^2 \vec{E} = -\epsilon_0 \mu_0 \frac{\partial^2 \vec{E}}{\partial t^2} - \mu_0 \frac{\partial^2 \vec{P}}{\partial t^2}. \quad (2.92)$$

This equation describes the wave equation in a general form. From this equation, we can see that the time varying polarization is a source term for new components of fields. To clearly see the influence of the nonlinear polarization, we will split it into two parts, the linear and the nonlinear polarization [51]:

$$\vec{P} = \vec{P}^{(L)} + \vec{P}^{(NL)}. \quad (2.93)$$

Using this splitting also in the relation between the electric field yields:

$$\vec{D} = \epsilon_0 \vec{E} + \vec{P}^{(L)} + \vec{P}^{(NL)} = \vec{D}^{(L)} + \vec{P}^{(NL)}. \quad (2.94)$$

Inserting this into Equation 2.92 leads to [51]:

$$-\nabla^2 \vec{E} + \mu_0 \frac{\partial^2 \vec{D}^{(L)}}{\partial t^2} = -\mu_0 \frac{\partial^2 \vec{P}^{(NL)}}{\partial t^2}. \quad (2.95)$$

In the case of a dispersive medium, we have to take the frequency components into account. Thus, we would have to split the linear polarization in the following manner [51]:

$$\vec{D}^{(L)} = \epsilon_0 \vec{E} + \vec{P}^{(L)} \quad (2.96)$$

$$= \frac{\epsilon_0}{2\pi} \int E(\omega) e^{-i\omega t} d\omega + \frac{\epsilon_0}{2\pi} \int \chi^{(1)}(\omega) E(\omega) e^{-i\omega t} d\omega \quad (2.97)$$

$$= \frac{\epsilon_0}{2\pi} \int \left(1 + \chi^{(1)}(\omega)\right) E(\omega) e^{-i\omega t} d\omega \quad (2.98)$$

$$= \frac{\epsilon_0}{2\pi} \int \epsilon^{(1)}(\omega) E(\omega) e^{-i\omega t} d\omega. \quad (2.99)$$

Furthermore, we represent the electric, linear displacement and polarization fields as the sum of their various frequency components [51]:

$$\vec{E}(r, t) = \sum_n \vec{E}_n(r, t) \quad (2.100)$$

$$\vec{P}^{(NL)} = \sum_n \vec{P}_n^{(NL)}(r, t). \quad (2.101)$$

This leads us to [51]:

$$-\nabla^2 E_n - \epsilon_0 \mu_0 \epsilon^{(1)}(\omega_n) \frac{\partial^2 E_n}{\partial t^2} = -\mu_0 \frac{\partial^2 P_n^{(NL)}}{\partial t^2}. \quad (2.102)$$

This means, we have to fulfill this equation for each frequency component. With this equation we can now calculate how the waves behave when propagating through a nonlinear waveguide which we will do next for the case of SFG.

Sum Frequency Generation

For this we assume two input fields with a frequency ω_1 and ω_2 which generate a field with frequency $\omega_3 = \omega_1 + \omega_2$. This condition is called energy conservation and must be fulfilled for every process. To find a solution, we assume a monochromatic guided mode propagating in z-direction through our waveguide of the form [54]:

$$E_i(x, y, z, t) = A_i(z)\mathcal{E}_i(x, y)e^{i(\beta_i z - \omega_i t)} + cc. \quad (2.103)$$

with $A_i(z)$ as the amplitude along the propagation direction, $\mathcal{E}_i(x, y)$ the spatial mode distribution and $\beta_i = \frac{n_i \omega_i}{c}$ as the propagation constant. We insert this ansatz into the nonlinear wave equation (see Equation 2.102) and take the following points into consideration when calculating the derivatives. First, E_3 solves the linear wave equation. Second, we assume that A_3 changes only slowly with z which is called the slowly varying amplitude approximation [51]:

$$\left\| \frac{\partial^2 A_3}{\partial z^2} \right\| \ll \beta_3 \left\| \frac{\partial A_3}{\partial z} \right\|. \quad (2.104)$$

Third, we neglect the complex conjugate terms without losing equality. This results in:

$$2i\beta_3 \frac{\partial A_3(z)}{\partial z} \mathcal{E}_3(x, y)e^{i(\beta_3 z - \omega_3 t)} = \mu_0 \frac{\partial^2 P_3(x, y, z, t)}{\partial t^2}. \quad (2.105)$$

Next, we use the relation between the nonlinear polarization and the E-field (see Equation 2.86) and insert the ansatz for our E-field. Combining this relation with Equation 2.105 and calculating the derivatives leads to:

$$\mathcal{E}_3(x, y) \frac{\partial A_3(z)}{\partial z} = \frac{i\mu_0 \epsilon_0 d\omega_3^2}{\beta_3} A_1(z)A_2(z)\mathcal{E}_1(x, y)\mathcal{E}_2(x, y)e^{i\Delta\beta z} \quad (2.106)$$

with $\Delta\beta = \beta_1 + \beta_2 - \beta_3$ as the phase mismatch. Lastly, we multiply the equation by the field distribution $\mathcal{E}_3(x, y)$ and integrate over the area assuming orthogonality of the modes which yields:

$$\frac{\partial A_3(z)}{\partial z} = \frac{i\mu_0 \epsilon_0 d\omega_3}{n_3} \theta A_1(z)A_2(z)e^{i\Delta\beta z}. \quad (2.107)$$

Here, we define θ as the mode overlap given by:

$$\theta = \int \int \mathcal{E}_1(x, y)\mathcal{E}_2(x, y)\mathcal{E}_3^*(x, y)dx dy. \quad (2.108)$$

Similar calculations for E_1 and E_2 result in a set of three coupled equations [54]:

$$\frac{\partial A_1(z)}{\partial z} = \frac{i\mu_0 \epsilon_0 d\omega_1}{n_1} \theta A_2(z)^* A_3(z)e^{i\Delta\beta z} \quad (2.109)$$

$$\frac{\partial A_2(z)}{\partial z} = \frac{i\mu_0 \epsilon_0 d\omega_2}{n_2} \theta A_1(z)^* A_3(z)e^{i\Delta\beta z} \quad (2.110)$$

$$\frac{\partial A_3(z)}{\partial z} = \frac{i\mu_0 \epsilon_0 d\omega_3}{n_3} \theta A_1(z)A_2(z)e^{i\Delta\beta z}. \quad (2.111)$$

With these three equations we can calculate the behavior of each field in our second order nonlinear waveguide.

The three equations depend on the phase mismatch $\Delta\beta$. Thus, next, we investigate how the amplitude A_3 changes with the propagation length in relation to the phase mismatch.

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Phase matching

First, we assume that we only have a small conversion efficiency of the input fields which means that the amplitudes A_1 and A_2 are constant [51]. For the special case $\Delta\beta = 0$, the spatial variation of A_3 is constant [51]:

$$\frac{\partial A_3(z)}{\partial z} = \text{const.} \quad (2.112)$$

This leads to the fact that the amplitude A_3 increases linearly with z and the intensity quadratically. This special case ($\Delta\beta = 0$) is called perfect phase matching (see Figure 2.13). This condition leads to a fixed phase relation between the generated wave and the nonlinear material polarization. Moreover, with this condition the highest energy transfer from A_1 and A_2 to A_3 is possible. However, perfect phase matching is not always fulfilled for arbitrary wavelengths combinations due to the dispersion of materials. [51] This means that the refractive index of a material depends on the frequency leading to the fact that the phase mismatch is frequency dependent:

$$\Delta\beta = \beta_1 + \beta_2 - \beta_3 = \frac{\omega_1 n_1(\omega_1)}{c} + \frac{\omega_2 n_2(\omega_2)}{c} - \frac{\omega_3 n_3(\omega_3)}{c}. \quad (2.113)$$

Thus, in a next step it is important to discuss the case $\Delta\beta \neq 0$.

If $\Delta\beta \neq 0$, the derivative of A_3 is not constant anymore, but we have to integrate over the length L of the nonlinear medium which results in [51]:

$$A_3(L) = \frac{i\mu_0\epsilon_0 d c \omega_3}{n_3} \cdot \theta \cdot A_1 \cdot A_2 \left(\frac{e^{i\Delta\beta L} - 1}{i\Delta\beta} \right). \quad (2.114)$$

To get the power, we use the following relation between the power and the amplitude $|A_i|^2 = \frac{2}{\epsilon_0 c n_i} P_i$ given by the time averaged Poynting vector. With this relation we obtain [51]:

$$P_3 = \frac{2d^2\omega_3^2}{c^3\epsilon_0 n_1 n_2 n_3} \cdot \theta^2 \cdot P_1 \cdot P_2 \cdot L^2 \cdot \text{sinc}\left(\frac{\Delta\beta L}{2}\right)^2. \quad (2.115)$$

Thus, the power of the generated wave depends sinc²-like on $\Delta\beta$. If $\Delta\beta = 0$, we get the highest power of the generated field. If $\Delta\beta \neq 0$ the power decreases significantly which is also shown in Figure 2.12 in which we plot the power of the generated wave as a function of $\Delta\beta L/2$. Moreover, the power of the generated field depends on the d-tensor element and the length, but also on the mode overlap of the interacting fields. This means, to achieve a higher power, we have to either use a higher d-tensor element, increase the length or we have to boost the mode overlap.

To achieve $\Delta\beta = 0$ which is also called momentum conservation, different approaches can be used such as angle tuning or birefringent phase matching. Using these approaches leads to the fact that the phase matching condition is only fulfilled for specific wavelength combinations. To achieve phase matching at various wavelengths, in this thesis, we will use quasi phase matching [55].

Quasi phase matching

In the method of quasi phase matching, the nonlinear properties of the crystal are altered periodically to generate a grating vector with which momentum conservation can be achieved

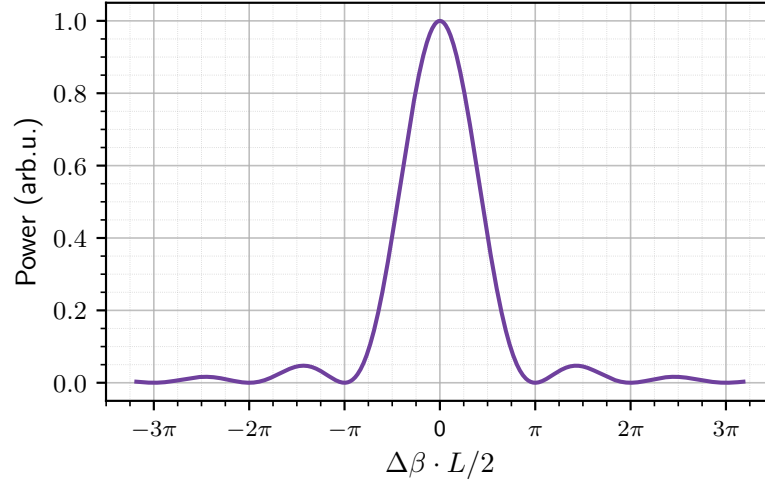


Fig. 2.12: Power as a function of the phase mismatch and length. Only if the phase mismatch is zero, the highest power can be reached. If $\Delta\beta \neq 0$ the power decreases.

for arbitrary wavelength combinations. To understand this process, we will look again at Equation 2.115. Here, we can see that the power of the generated field depends on the length of the crystal as well as on the phase mismatch. Now, we can see that the power P_3 is the highest if the sinus is equal to 1. This is the case if the length L is equal to $L_{\text{coh}} = \frac{\pi}{\Delta\beta}$. This length is also called coherence buildup length up until which the power P_3 will increase [51, 55]. After this length, the power of the generated wave P_3 will decrease again. We can also explain this via interference of the interacting waves. If we assume for example SHG, we have an input field which generates an output field along the waveguide. The input field and the generated output field will travel through the waveguide with different velocities due to the different refractive indices if we assume a birefringent crystal. This means the two waves have a changing phase relation and the coherence buildup length is the length at which the two waves have a phase difference of π . The phase of the input field determines the phase of the locally generated output fields. Thus, the locally generated output fields will have a phase difference which means they can either interfere constructively or destructively. To build up the power of the generated output field, we use quasi phase matching which benefits the constructive interference between the locally generated output fields. For this, we invert the nonlinear polarization of the crystal each coherence buildup length which introduces a grating vector $\frac{2\pi}{\Lambda}$. This alters the momentum conservation condition as follows [51]:

$$\Delta\beta = \beta_1 + \beta_2 - \beta_3 - \frac{2\pi}{\Lambda} \approx 0. \quad (2.116)$$

Here, Λ is the so called poling period which is defined as [51]:

$$\Lambda = 2L_{\text{coh}} = \frac{2\pi}{\Delta\beta}. \quad (2.117)$$

This leads to a change in the sign of d_{eff} and thus in a change in the sign of $\Delta\beta$. This means, the phase difference is flipped and the energy will not flow back to the input wave after the coherence length, but will further flow into the output wave and increase its power (see Figure 2.13). [55] The phase matching function is often shown in a 2D

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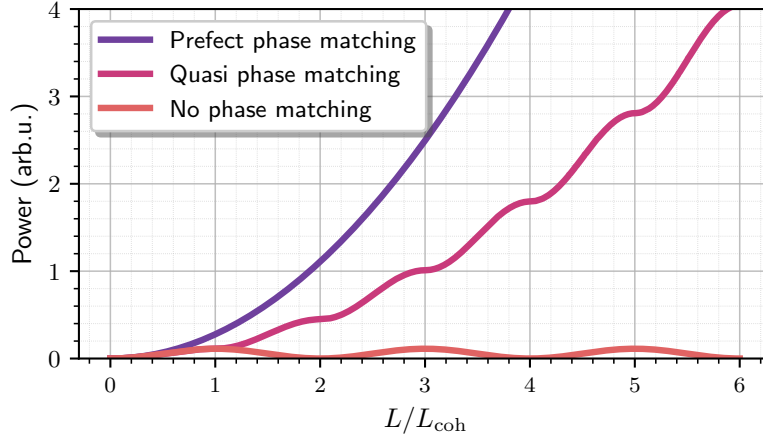


Fig. 2.13: Power of the generated wave as a function of the length. For no phase matching between the interacting fields, the power only increases up to once the coherence buildup length. Afterwards it decreases again. If we use quasi phase matching the power increases further than the coherence buildup length since we invert the nonlinear polarization each coherence buildup length. For perfect phase matching the power increases quadratically. Redrawn after [55].

plot. In Figure 2.14, we show the phase matching function of a type 0 SFG process at 1550 nm and 810 nm to generate light at 532 nm in a TFLN waveguide using quasi phase matching. The phase matching function is shown as a function of the input and output wavelengths. We can see the sinc^2 -like structure and the light can be generated where the phase matching is fulfilled. The shape of the phase matching function generally depends on the material dispersion as it is the case for titanium indiffused waveguides in lithium niobate. But, as we will later see, for TFLN waveguides the dispersion is not solely determined by the material but is strongly influenced by the waveguide geometry. This control over the dispersion is a clear advantage compared to titanium indiffused waveguides

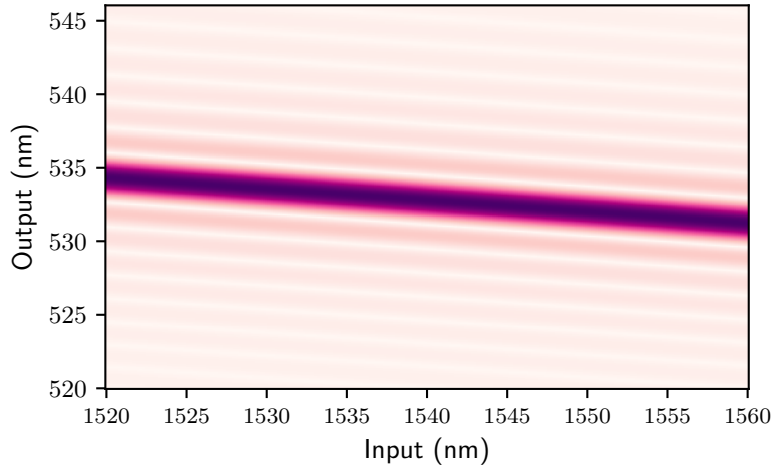


Fig. 2.14: Phase matching function of a type 0 SFG process. In this process light at 1550 nm and 810 nm generate light at 532 nm. We can see the sinc^2 -shape of the phase matching. The light can only be generated in those regions.

2.4. Thin-Film Lithium Niobate

since we can achieve nonlinear processes with engineered dispersion properties tailored to specific applications. This method called dispersion engineering will be discussed in detail later. The physical implementation of quasi phase matching will also be explained later.

Now, we understand how a second-order susceptibility leads to the conversion of light to a different frequency. Next, we want to introduce the material platform which we will use in this thesis to realize nonlinear processes.

2.4 Thin-Film Lithium Niobate

Lithium niobate (LN) was the material of choice for integrated (quantum) optics for the last decades. However, devices based on standard LN face the drawback of a low mode confinement due to the low index contrast difference of the waveguides. Thus, the devices are relatively large compared to circuits based on silicon photonics. Therefore, the devices are bulky and less efficient. But, LN was revolutionized with the possibility of manufacturing thin films of LN bonded on an insulator layer (called lithium niobate on insulator (LNOI) or thin-film lithium niobate (TFLN)) [9] and realizing nano-structured waveguides [10, 11, 12, 13, 14]. Due to this, devices with reduced footprints and enhanced efficiencies became feasible. Thus, today, TFLN is an actively researched material platform exploring highly efficient and compact devices and circuits for integrated (quantum) optics applications [16, 56]. Hence, TFLN is also the material of choice for this thesis. Therefore, in this chapter, we will introduce the material properties of LN and the layer system. Afterwards, we will discuss waveguides in TFLN and how those are modeled. Lastly, we will introduce the nonlinear properties of TFLN and discuss how nonlinear processes can be realized.

2.4.1 Material Properties

Lithium niobate is a artificial crystal that consists of lithium, niobium and oxygen atoms [57, 58]. It belongs to the trigonal crystal system. LN is a uniaxial crystal with a birefringence that is defined by the ordinary n_o and extra-ordinary refractive index n_e . Moreover, LN is a non-centro symmetric crystal and thus exhibits a second-order nonlinearity as discussed before [57]. Due to this lack of inversion symmetry LN is a ferroelectric crystal exhibiting a spontaneous polarization. This polarization can be inverted by applying a strong electric field. This is later used for periodic poling of the material. Furthermore, LN has a large transparency window and large electro-optical coefficients [57, 58, 16].

All of these parameters make LN an ideal platform for integrated photonics. The waveguides in standard LN are fabricated mostly via a titanium indiffusion or proton exchange. This leads to a small refractive index contrast which makes devices on standard LN relatively large compared to other photonic platforms [16, 5]. To miniaturize the devices TFLN can be utilized since it combines the outstanding optical properties of LN with a high index contrast as discussed next.

Thin-film lithium niobate consists of a sub-micron layer of LN bonded onto an insulator layer which is connected to a handle substrate (see Figure 2.15). In this thesis, we will use a thin-film thickness of LN of 600 nm bonded on to a silicon dioxide insulator layer in a x-cut and y-propagation configuration. For linear devices (such as directional couplers and racetrack resonators) we use congruent LN, but for nonlinear devices (all devices with periodic poling) we use LN with a 5 mol% magnesium dioxide doping because of the lower

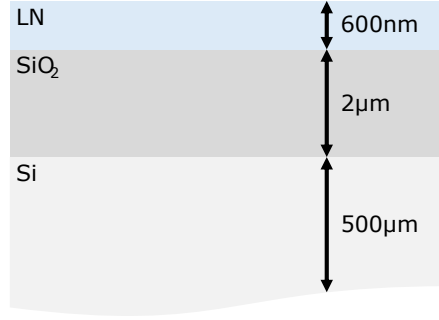


Fig. 2.15: The layer system of TFLN consists of a thin layer of LN bonded to an insulator layer, here silicon dioxide. Beneath the SiO_2 a handle substrate is bonded.

coercive field strength which is important for periodically poling of the LN [59, 60]. The handle substrate is silicon or LN.

The light is guided in the LN thin-film and thus, TFLN inherits the optical properties from bulk LN. However, TFLN also has a much higher refractive index contrast compared to standard LN waveguides. For example, titanium indiffused waveguides in LN have an index contrast of around 0.001 to 0.04 [16]. TFLN has a index contrast of up to two order of magnitude higher due to the SiO_2 underneath the LN. This allows for a high vertical confinement [16]. To fully exploit the higher confinement also in the horizontal directional waveguides are fabricated in the LN film. Next, we will discuss the shape as well as the modeling of waveguides in TFLN.

2.4.2 Modeling and Nonlinear Processes of Waveguides in TFLN

To take full advantages of the higher confinement of TFLN waveguide structures are used. In this thesis, ridge waveguides are fabricated as shown Figure 2.16. The waveguide is defined by the thin-film thickness t , the etching depth h , the top width w , the angle a and the SiO_2 cladding thickness c . The top width of TFLN waveguides are around $1\mu\text{m}$ to

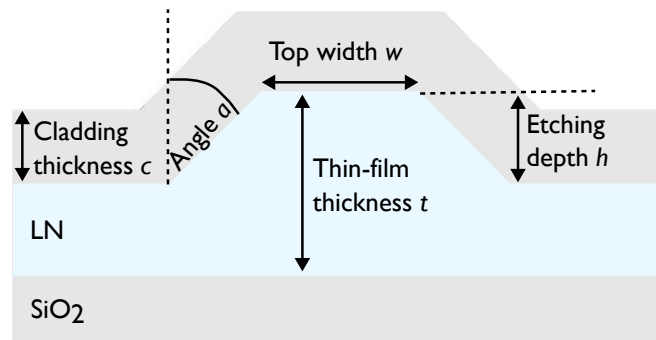


Fig. 2.16: Schematic of an etched waveguide in TFLN. The waveguide is defined by the geometry parameters which are the thin-film thickness t , the etching depth h , the top width w , the waveguide angle a and the cladding thickness c .

2.4. Thin-Film Lithium Niobate

$3\text{ }\mu\text{m}$ depending on the number of guided modes and thus the mode size area is around $1\text{ }\mu\text{m}^2$. The light is guided in the LN rib and thus the TFLN waveguides have the same optical properties as waveguides in standard LN. But, now the refractive index contrast is different compared to standard LN due to the SiO_2 underneath the LN thin-film. This leads to a refractive index contrast of around 0.5 which is two orders of magnitude higher than of standard LN waveguides [16]. This decreases significantly the bending radii to a few ten micrometer compared to the millimeter-range bending radii in standard LN enabling compact devices [16]. To design those devices it is crucial to model the dispersion of those waveguides to predict their functionality. Moreover, this is also important for dispersion engineering of quantum devices. Thus, we will discuss next how to model TFLN waveguides.

Modeling of Waveguides using *Lumerical*

To model the dispersion of TFLN waveguides, we have to solve the vectorial wave equation, see Equation 2.13. This is a partial differential equation (PDE) which can in our case not be solved analytically. To solve this numerically we use the Finite Difference Method (FDM). The FDM is a way to numerically solve PDE [61]. It can be applied to different kinds of problems with different boundary shapes or boundary conditions and also for regions with more than one material [62]. In the FDM the partial differential equation is solved only at finite points by evaluating difference quotients [62]. This solving method is implemented in the *MODE* solver of *Lumerical* which is a finite difference eigenmode solver. The waveguide model as well as the simulation settings were already implemented during the PhD thesis of Dr. Matteo Santandrea [63]. An image of the implemented waveguide model is shown in Figure 2.17. The model consists of a $10\text{ }\mu\text{m}$ thick SiO_2 layer which represents the insulator layer. On top of this is the LN thin-film of which the thickness can be changed. Furthermore, we can define an etching depth, a top width and the angle of the waveguide. Moreover, we can apply a cladding layer. The handle substrate is not implemented since the mode will not reach into it. The mesh of the simulation has a span of $6\text{ }\mu\text{m} \times 10\text{ }\mu\text{m}$ with a size of the

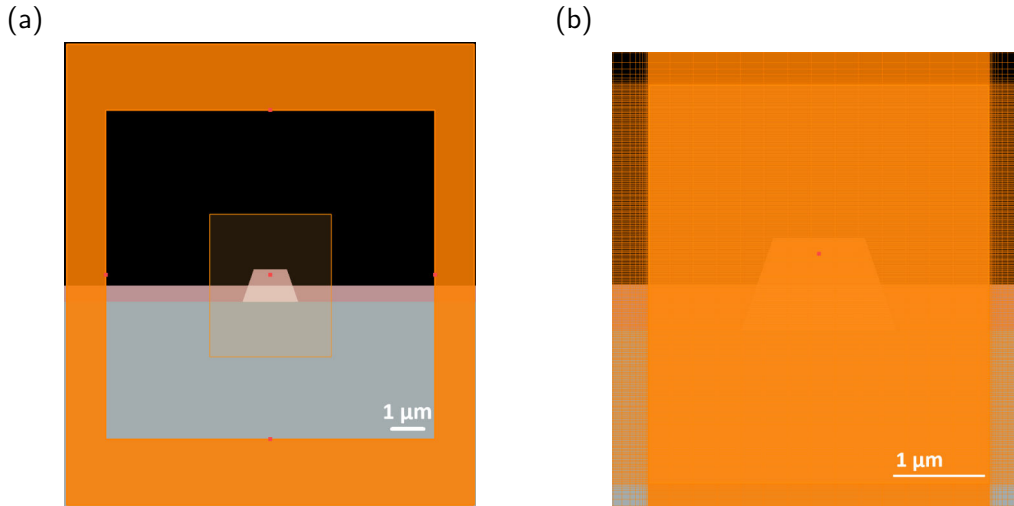


Fig. 2.17: Images from the simulation software *Lumerical*. (a) The waveguide implemented in *Lumerical* without a cladding layer is shown. Furthermore, we can see a big rectangular which is the mesh region. The smaller rectangular around the waveguide is the refined mesh which is shown in (b).

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mesh grid of $0.1\ \mu\text{m}$. *Lumerical* provides the opportunity to apply a refined mesh for more crucial regions which is in our case the waveguide. Thus, we apply a refined mesh with a mesh size of $0.01\ \mu\text{m} \times 0.005\ \mu\text{m}$ with a span of $1\ \mu\text{m}$ from the edge of the waveguide. As boundary conditions, we use perfectly matched layers (PML) which means the light outside the waveguide is absorbed which minimizes reflections. The model and settings were optimized by Dr. Matteo Santandrea.

To simulate the dispersion of TFLN waveguides, we simulate the effective refractive index between $400\ \text{nm}$ and $2500\ \text{nm}$ in 22 wavelength steps and apply a fit based on LN Sellmeier equation. As refractive index of the lithium niobate thin-film, we use for congruent LN the Sellmeier equations reported by Jundt [64] and Edwards et al. [65] and for 5 mol% MgO doped LN the Sellmeier equations reported by Gayer et al. [66]. We use for the refractive index of the silicon dioxide the values provided by *Lumerical*. These simulations are the basis for realizing any linear, nonlinear or quantum device in TFLN. To achieve nonlinear processes, we need to know which processes are allowed by the d-Tensor.

Nonlinear Processes in TFLN Waveguides

To realize an efficient nonlinear process in TFLN, we use periodic poling to achieve quasi phase matching. Therefore, we next discuss which tensor element LN has and which we can use for our x-cut y-propagation samples. Afterwards, we will describe the fabrication of the periodic poling.

Lithium niobate has the following d-Tensor [58]:

$$d_{\text{LN}} = \begin{bmatrix} 0 & 0 & 0 & 0 & d_{31} & -d_{22} \\ -d_{22} & d_{22} & 0 & d_{31} & 0 & 0 \\ d_{31} & d_{31} & d_{33} & 0 & 0 & 0 \end{bmatrix} \quad (2.118)$$

Inserting this into Equation 2.86 and taking into account that we have y-propagation meaning that we have no E-field in y-direction yields the following possible processes:

- type 0: $\text{TE} \rightarrow \text{TE} + \text{TE}$
- type 1: $\text{TE} \rightarrow \text{TM} + \text{TM}$
- type 2: $\text{TM} \rightarrow \text{TM} + \text{TE}$

Thus, we can realize these three different processes in our x-cut and y-propagation samples. To realize them on a sample, we first have to periodically pole the sample and fabricate waveguides which is explained next.

Fabrication of Periodically Poled Waveguides

The fabrication of (periodically poled) waveguides in TFLN is a challenging task. Thus, throughout the scope of this thesis, the fabrication process was improved by other group members which means that each sample used in this thesis can be fabricated in a slightly different way. Here, we only present the most recent process since the focus of this thesis is not the fabrication and all samples were fabricated by other very kind group members.

The LN thin-film is poled before etching the waveguides. To pole the thin-film, we first structure finger electrodes on the sample surface using an electron-beam lithography. After a metal deposition, a high-voltage pulse is applied to invert the spontaneous polarization in the crystal. More information about mastering poling of TFLN can be found in [67].

After poling, the waveguides in TFLN are fabricated using an electron-beam lithography and physical dry etching. If required, we can deposit a SiO₂ cladding layer. In the final step, the end faces of the sample are chemomechanically polished to improve the free-space coupling efficiency into the waveguides.

This is the material platform which we investigate in detail as a material for integrated quantum optics. Thus, we will introduce the quantum phenomena utilized in this thesis in the next chapter.

2.5 Quantum Optics

In this chapter, we discuss the relevant quantum optical principles and processes that are used in this thesis. We start with discussing parametric down-conversion as a process for generating single photons and discuss correlations between the generated photons. Afterwards, we explain the Hong-Ou-Mandel interference as the most basic quantum optical phenomenon. Lastly, we briefly introduce the quantum pulse gate which is based on a dispersion-engineered SFG process.

2.5.1 Parametric Down-Conversion

Parametric down-conversion (PDC) is a nonlinear quantum process in which a mother photon spontaneously decays into two daughter photons, called signal and idler, inside a second order nonlinear crystal [68]. This processes is similar to the nonlinear process discussed above and thus we can use the same nomenclature. Hence, for this process to take place, we also have to fulfill energy conservation

$$\hbar\omega_p = \hbar\omega_s + \hbar\omega_i \quad (2.119)$$

so that the frequencies of the signal ω_s and idler photon ω_i add up to the pump frequency ω_p . Moreover, for an efficient process we again need to fulfill momentum conservation

$$\beta_p - \beta_i - \beta_s - \frac{2\pi}{\Lambda} = 0 \quad (2.120)$$

meaning that the propagation constants of the signal β_s and idler photon β_i add up to the propagation constant of the pump photon β_p which we achieve with quasi phase matching and the poling period Λ as an additional grating vector.

To describe this process in a quantum fashion, we need a Hamiltonian of the system. It can be shown that the Hamiltonian of a system with a classical pump field and a single spatial mode operation in the waveguide is given by [69]:

$$\hat{H}_{PDC} = B \cdot \int d\omega_s d\omega_i f(\omega_i, \omega_s) \hat{a}^\dagger(\omega_s) \hat{b}^\dagger(\omega_i). \quad (2.121)$$

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Here, B combines all constants and is determined by the field overlap of the interacting modes and the pump power, $\hat{a}^\dagger$ and $\hat{b}^\dagger$ are the creation operators for the signal and idler photon and the function $f(\omega_i, \omega_s)$ is the so called joint spectral amplitude (JSA) which gives information about the spectral and temporal properties of the generated signal and idler photons. The JSA is simultaneously given by the energy conservation (pump function) and momentum conservation (phase matching function) [70, 71]:

$$f(\omega_i, \omega_s) = \alpha(\omega_s + \omega_i) \cdot \phi(\omega_s, \omega_i). \quad (2.122)$$

Here, for simplicity we assume a Gaussian pump for the pump function $\alpha(\omega_s + \omega_i)$:

$$\alpha(\omega_s + \omega_i) = e^{-\frac{(\omega_s + \omega_i)^2}{2\sigma}}. \quad (2.123)$$

The phase matching function $\phi(\omega_s, \omega_i)$ is defined as [71]:

$$\phi(\omega_s, \omega_i) = \text{sinc}\left(\frac{\Delta\beta L}{2}\right) e^{-\frac{i\Delta\beta L}{2}} \quad (2.124)$$

which has the same form as derived in Equation 2.115. To visualize the spectral properties and correlations, this is usually shown in the signal and idler plane which we discuss next.

Joint Spectral Amplitude

In Figure 2.18, we can see three 2D plots that represent the phase matching function, the pump function and the joint spectral intensity (JSI) which is the absolute square of the JSA shown as a function of the signal and idler frequencies.

Generally, the phase matching function is determined by the dispersion of the material. TFLN, however, has the advantage that the dispersion does not only depend on the material dispersion, but it strongly depends on the waveguide geometry due to the high mode confinement. This opens up the possibility to tailor the dispersion on a broader scope compared to using titanium indiffused waveguides in lithium niobate to realize PDC sources optimized for various applications. Thus, by engineering the dispersion, we can tune the group velocities and therefore the phase matching function, as we will later discuss in chapter 5.

An important property for the classification of the phase matching which also provides information about the correlations of the generated photons is the phase matching angle. For a PDC process the phase matching angle is defined as [72]:

$$\alpha_{\text{PDC}} = -\arctan\left(\frac{v_{g,\text{pump}}^{-1} - v_{g,\text{signal}}^{-1}}{v_{g,\text{pump}}^{-1} - v_{g,\text{idler}}^{-1}}\right). \quad (2.125)$$

Here, v_g describes the group velocity. This means, to tailor the phase matching angle, we need to have precise control over the group velocities. TFLN provides this control due to the waveguide-dependent dispersion.

Depending on the angle and thus, on the relation of the group velocities of the pump, signal and idler photon, the phase matching is categorized. For this thesis, we discuss three different cases. All three cases are shown in Figure 2.18 where the phase matching function as well as the pump function and resulting JSI are displayed in the signal and idler plane.

In Figure 2.18 (a), we can see a non-engineered phase matching function for which the phase

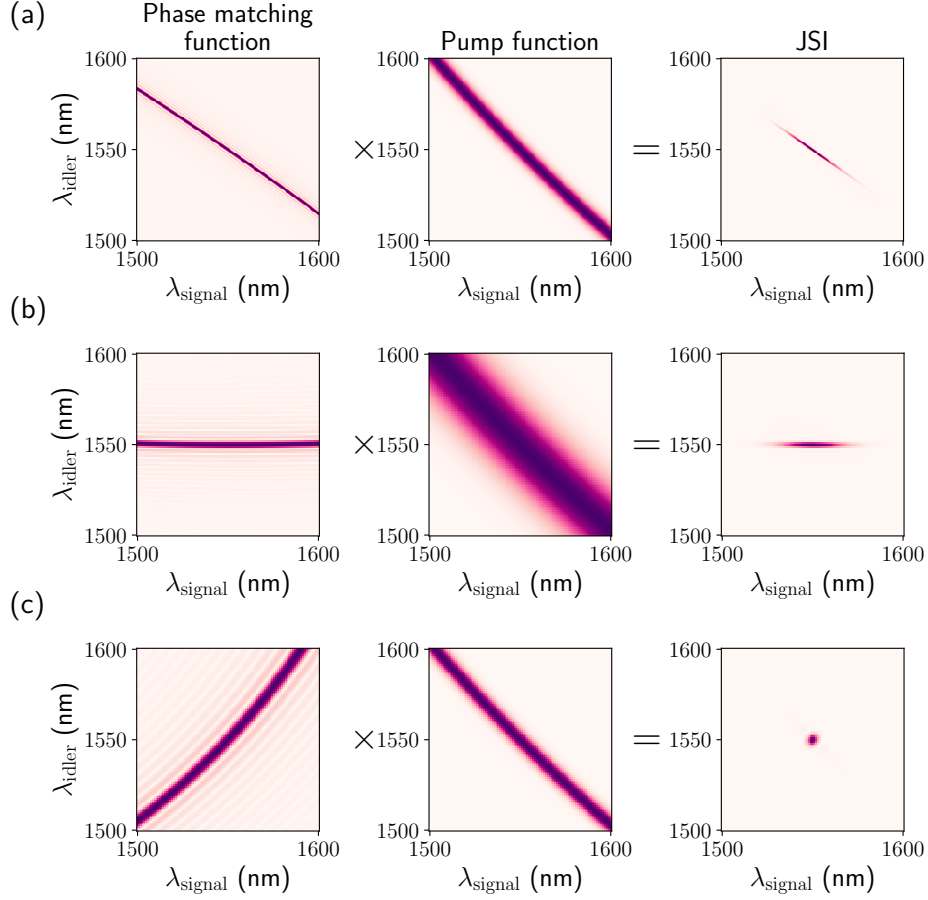


Fig. 2.18: JSI of different non-engineered and engineered phase matching functions. (a) The first plot shows a non-engineered phase matching function with a negative phase matching angle resulting in correlations in the JSI. (b) Asymmetric group velocity matching is shown which leads to no correlations if the phase matching and pump bandwidth match or the pump bandwidth is larger than the phase matching bandwidth. (c) A phase matching function with a angle of 45° is shown which leads to no correlations in the JSI if the bandwidth of the phase matching and pump are matched.

matching angle is negative. If we multiply this phase matching function with the pump function which is always under an angle of -45° due to energy conservation, the resulting JSI shows a ellipse under -45° . Therefore, the signal and idler photon are correlated since knowing the signal wavelength leads to knowing the idler wavelength.

A challenge is to generate uncorrelated photons. To achieve a uncorrelated JSI, a phase matching angle between 0° and 90° is necessary as well as a matched pump bandwidth [72]. Two special cases are if $\alpha_{PDC} = 0^\circ$ and if $\alpha_{PDC} = 45^\circ$. To reach these special angles, dispersion engineering is often used which will be explained in section 5.1. We will discuss each case individually next.

If the phase matching function is parallel to the x- or y-axis, it is called asymmetric group velocity matching [72, 73]. To achieve this, one of the generated photons travels with

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the same group velocity as the pump photon. If the signal/idler photon has the same group velocity as the pump photon, we can see that Equation 2.125 gives a values of $\alpha_{PDC} = 0^\circ/90^\circ$. The correlation in the output state depends on the pump bandwidth. If the pump bandwidth is much narrower than the phase matching bandwidth, the output state is an ellipse under -45° and thus the JSI is correlated. If we increase the pump bandwidth larger than the phase matching bandwidth, we get an ellipse along the x-axis/y-axis which provides an uncorrelated JSI. This is shown in Figure 2.18 (b).

The other important case is called symmetric group velocity matching [74]. Here, the group velocity of the generated signal and idler photons are symmetric around the group velocity of the pump photon. This yields in Equation 2.125 an angle of 45° . To achieve an uncorrelated state, we need to match the pump bandwidth to the phase matching bandwidth. If these two are the same, we achieve a circle in the JSI and thus, there are no correlations between the signal and idler photons as it can be seen in Figure 2.18 (c).

A more mathematically approach towards correlations in the JSI can be made via the Schmidt decomposition which will be discussed next.

Schmidt Decomposition

Any 2D function has in general correlations between its variables. With the Schmidt decomposition, one can find a special basis called the Schmidt basis which allows to represent the correlations of the function via the Schmidt coefficients.

A 2D function can be described as the superposition of two sets of orthonormal functions [70, 71]:

$$F(\omega_s, \omega_i) = \sum_j \sqrt{\lambda_j} g_j(\omega_s) h_j(\omega_i). \quad (2.126)$$

Here, $\{g_j(\omega_s)\}$ and $\{h_j(\omega_i)\}$ are sets of orthonormal functions also called Schmidt basis and $\sqrt{\lambda_j}$ are non-negative real numbers which have to satisfy $\sum_j \lambda_j = 1$ named Schmidt coefficients [75]. The Schmidt decomposition of a function that depends on two parameters such as the JSA describes how many orthonormal functions are necessary to describe the function. In other words, we find the minimal number of orthonormal functions to describe F .

The Schmidt coefficients and modes are a useful tool to analyze correlations in the JSA. If the JSA can be described by one Schmidt coefficient, the state is separable and therefore the JSA is uncorrelated. This is the case if the JSA is a circle or an ellipse along one axis since then we can use one pair of Schmidt modes to describe it as for example if $\alpha_{PDC} = 0^\circ$ or $\alpha_{PDC} = 45^\circ$.

If we have correlations in the JSA by having for example a negative phase matching angle, we need more than one Schmidt coefficient to describe the state. Thus, the state is described by more than one pair of Schmidt modes and therefore not separable meaning we will measure correlations.

This is a theoretical way to characterize a PDC source. Next, we will discuss experimentally figures of merit for the characterization of a PDC source.

Characterization of a PDC Source

A PDC source can be experimentally characterized by different figures of merit. In this thesis, we will focus on three different ones namely the Klyshko efficiency, the heralded



Fig. 2.19: Measurement scheme to evaluate the Klyshko efficiency.

second-order correlation function $g^2(0)_h$ as well as the brightness of the source. The first important parameter to discuss is the Klyshko efficiency [76] which describes the efficiency of the photons from the source to the detection. It evaluates how likely it is to measure an idler photon when a signal photon has been detected or vice versa. The measurement scheme to measure the Klyshko efficiency is shown in Figure 2.19. We pump our PDC source and measure the generated signal and idler photons with detectors. The PDC generates a number of pairs per second which is called the PDC rate R_{PDC} . The rate of measuring a signal photon R_s or idler photon R_i is given by the efficiency to measure the photon η_s/η_i times the PDC rate R_{PDC} :

$$R_s = \eta_s \cdot R_{\text{PDC}} \quad (2.127)$$

$$R_i = \eta_i \cdot R_{\text{PDC}}. \quad (2.128)$$

Moreover, we can also measure the coincidences between the two arms and determine a coincidence rate $R_{s,i}$. This rate is given by the efficiency of detecting a signal photon η_s times the efficiency of detecting an idler photon η_i times the PDC rate R_{PDC} :

$$R_{s,i} = \eta_s \cdot \eta_i \cdot R_{\text{PDC}}. \quad (2.129)$$

By dividing the coincidence rate by the signal or idler rate, we get the Klyshko efficiency for the signal η_s and idler photon η_i :

$$\eta_s = \frac{R_{s,i}}{R_i} \quad (2.130)$$

$$\eta_i = \frac{R_{s,i}}{R_s}. \quad (2.131)$$

This equation is equal to one when the same amount of coincidences and singles are measured. However, if more single events than coincidences are measured, the Klyshko efficiency is reduced. This means, we lost one photon from the photon pair. Therefore, the Klyshko efficiency is a metric for single photon loss. The losses can be due to different reasons such as losses in the sample, losses due to coupling or detection loss. Since the Klyshko efficiency describes all setup losses it is a highly important parameter because we can calculate from the measured rates the actual generated PDC rate inside the sample. However, in a PDC source we do not only generate single pairs, we can also generate multiple photon pairs. To quantify this, a Hanbury-Brown-Twiss setup is used as shown in Figure 2.20. Here, we measure the signal arm directly but introduce a 50:50 beam splitter in the idler arm and measure both outputs. If we now assume that the PDC source generates one photon pair, the signal photon would directly be measured. The idler photon would either be reflected or transmitted at the beam splitter. Thus, the idler photon can only be detected in one of the two output arms of the beam splitter. However, if multiple photon pairs are generated, it is

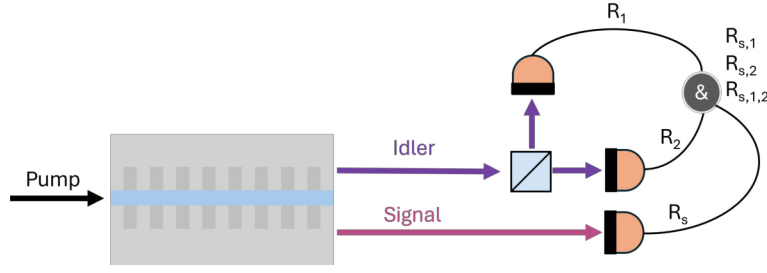


Fig. 2.20: Measurement scheme to measure the heralded second-order correlation function $g_h^{(2)}(0)$.

possible to measure a photon in both outputs of the beam splitter. This behavior can be evaluated by measuring the coincidences between all three outputs. This value is a central parameter of the heralded second-order correlation function $g_h^{(2)}(0)$ which is a figure of merit for the simultaneous generation of multiple photon pairs. To estimate the $g_h^{(2)}(0)$, we have to measure the single count rates of the signal arm R_s , the twofold coincidences between the signal arm and the idler 1 arm $R_{s,1}$ and the signal arm and the idler 2 arm $R_{s,2}$ as well as the threefold coincidence between all three outputs $R_{s,1,2}$. With this values we can calculate the heralded second-order correlation function $g_h^{(2)}(0)$ [69]:

$$g_h^{(2)}(0) = \frac{R_{s,1,2} \cdot R_s}{R_{s,1} \cdot R_{s,2}}. \quad (2.132)$$

For an ideal single photon source, the $g_h^{(2)}(0)$ is equal to zero since this means we did not measure any three-fold coincidences. If the PDC source generates multiple photon pairs, the $g_h^{(2)}(0)$ value increases.

The last important figure of merit is the brightness of the source. The brightness describes the performance of the source and gives information about how many pairs are generated per time, per pump power and per bandwidth. It is defined as the generated photon pair rate R_{PDC} over the pump power inside the waveguide P and the bandwidth $\Delta\omega$ of the generated photons:

$$B = \frac{R_{\text{PDC}}}{P \cdot \Delta\omega}. \quad (2.133)$$

A high brightness is favorable to achieve short measurement times and a high signal-to-noise ratio and low pump powers.

To conclude, a PDC process generates photon pairs. Depending on the dispersion and thus on the resulting phase matching function, the generated photons are correlated or decorrelated. There are different figures of merit to experimentally characterize a PDC. For a PDC source another important figure of merit can be measured via the Hong-Ou-Mandel interference which is discussed next.

2.5.2 Hong-Ou-Mandel Interference

The Hong-Ou-Mandel interference describes the interaction of two photons at a balanced, non-polarizing beam splitter [77]. If two indistinguishable photons enter one input port of

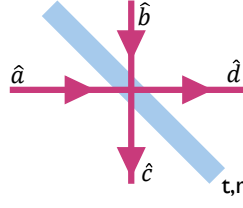


Fig. 2.21: Schematic of a beam splitter. The beam splitter operation is described by the reflectivity r and the transmissivity t . Two fields $\hat{a}$ and $\hat{b}$ enter the beam splitter and $\hat{c}$ and $\hat{d}$ are the output fields.

the beam splitter each, they bunch and leave the beam splitter together in one output port. To understand this behavior better, we first have a closer look at the quantum mechanical treatment of a beam splitter as shown in Figure 2.21. The two inputs are labeled $\hat{a}$ and $\hat{b}$ and the two outputs $\hat{c}$ and $\hat{d}$. Important for the description of a quantum mechanical beam splitter is that we also consider both input fields even if we only have one since an "empty" input is described by the vacuum field. The beam splitter has the reflectivity r and the transmissivity t . We can describe the process of the beam splitter with the beam splitter matrix [78]:

$$\begin{pmatrix} \hat{c} \\ \hat{d} \end{pmatrix} = \begin{pmatrix} r & t \\ t & r \end{pmatrix} \begin{pmatrix} \hat{a} \\ \hat{b} \end{pmatrix}. \quad (2.134)$$

By looking at the commutation relations, we obtain the following conditions for r and t [78]:

$$\|t\|^2 + \|r\|^2 = 1 \quad (2.135)$$

$$r \cdot t + r^* \cdot t^* = 0. \quad (2.136)$$

Consequently, for a asymmetric 50:50 beam splitter meaning that a π phase shift only occurs in one output arm the beam splitter matrix is as follows [78]:

$$\begin{pmatrix} \hat{c} \\ \hat{d} \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 1 \\ -1 & 1 \end{pmatrix} \begin{pmatrix} \hat{a} \\ \hat{b} \end{pmatrix}. \quad (2.137)$$

This leads to the following relation between the input and output creation operators [79]:

$$\hat{a}^\dagger = \frac{1}{\sqrt{2}} \cdot (\hat{c}^\dagger - \hat{d}^\dagger), \hat{b}^\dagger = \frac{1}{\sqrt{2}} \cdot (\hat{c}^\dagger + \hat{d}^\dagger). \quad (2.138)$$

If we now let these operators act on a vacuum state to calculate the output modes of the beam splitter [79]

$$|\Psi_{\text{out}}\rangle = \hat{a}^\dagger \hat{b}^\dagger |0, 0\rangle = \frac{1}{2} (\hat{c}^\dagger \hat{c}^\dagger + \hat{c}^\dagger \hat{d}^\dagger - \hat{c}^\dagger \hat{d}^\dagger - \hat{d}^\dagger \hat{d}^\dagger) |0, 0\rangle \quad (2.139)$$

we see that there are four different possibilities how two photons can interact on a beam splitter which are shown in Figure 2.22. The first possibility is that both photons are reflected at the beam splitter. Second, both photons can be transmitted. The third and fourth option are that one photon is transmitted and the other reflected such that both leave the beam

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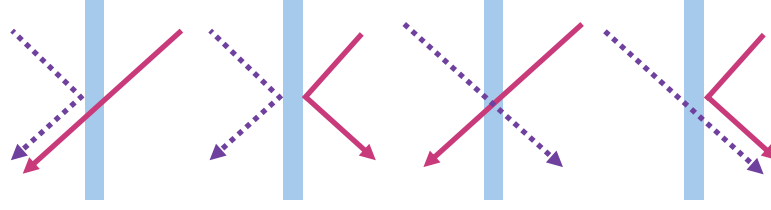


Fig. 2.22: Schematic of two photons impinging on a 50:50 beam splitter. There are four different outcomes, but if the two photons are indistinguishable they will always leave the beam splitter in one port.

splitter together. If the two photons are distinguishable, all four possibilities occur. However, if the photons are indistinguishable, the terms $\hat{c}^\dagger \hat{d}^\dagger - \hat{c}^\dagger \hat{d}^\dagger$ cancel out. This leaves us with:

$$|\Psi_{\text{out}}\rangle = \frac{1}{2} (\hat{c}^\dagger \hat{c}^\dagger - \hat{d}^\dagger \hat{d}^\dagger) |0, 0\rangle = \frac{1}{\sqrt{2}} (|2, 0\rangle - |0, 2\rangle). \quad (2.140)$$

This shows that the two photons always leave the beam splitter together. This effect is called photon bunching. It can be experimentally visualized by delaying one of the photons in time in relation to the second photon and measuring the coincidences between the two output ports of the beam splitter. If the two photons arrive at different times on the beam splitter, all four outcomes are possible, since the two photons are distinguishable due to the time delay. Consequently, we would measure coincidences between the two output ports. However, if the two photons arrive at the same time and have the same properties, photon bunching occurs and the two photons leave the beam splitter in the same output. Therefore, no coincidences are measured and thus the coincidences drop to zero (see Figure 2.23). This drop is called Hong-Ou-Mandel dip [77] and if the photons are truly indistinguishable

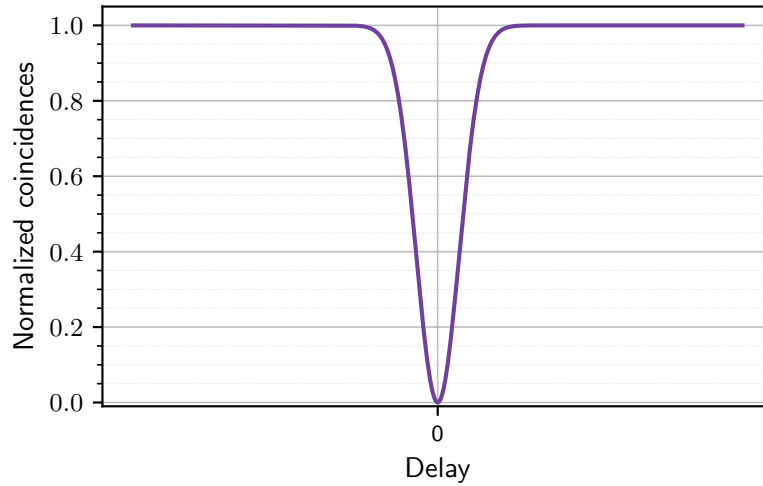


Fig. 2.23: The graph shows the Hong-Ou-Mandel dip by plotting the normalized coincidences of the two output ports of a beam splitter as a function of the time delay between two photons each entering one input port of the beam splitter. If the two photons share the same properties and arrive on the beam splitter at the same time, they leave the beam splitter in one output port. Consequently, no coincidences are measured and we see a dip.

the visibility of the dip is 100 %. If the two photons are not completely indistinguishable, the visibility of the dip is reduced. If we measure the HOM dip via a time delay, the width of the dip corresponds to the bandwidth of the photons. Thus, it is important that the bandwidth of the beam splitter is larger than the bandwidth of the photons. If this is not the case, the visibility is reduced. Furthermore, a 100 % visibility can only be achieved with a perfect 50:50 splitter. The maximal visibility of a non perfect splitter can be calculated via the relation [80]:

$$V_{\max} = \frac{2 \cdot \zeta \cdot (1 - \zeta)}{1 - 2 \cdot \zeta + 2 \cdot \zeta^2}, \quad (2.141)$$

where ζ is the reflectivity of the beam splitter.

Thus, photons from a PDC source with symmetric group velocity matching can be used to show HOM interference. Photons from a PDC source can also be used as an input for the quantum pulse gate which we will discuss next.

2.5.3 Quantum Pulse Gate

The quantum pulse gate (QPG) [81, 82, 83, 84, 85] is a temporal-mode selective frequency conversion device. It is based on a dispersion-engineered SFG process inside a periodically poled waveguide. A (quantum) input pulse propagates through the waveguide at the same group velocity as a bright pump pulse, whose complex spectrum is shaped in advance. The QPG is commonly used for the detection of temporal modes. Temporal modes describe wave-packet modes of quantum states of light. They are described by a complex amplitude $f(\omega)$, which combines the spectrum $\|f(\omega)\|^2$ and a spectral phase. Temporal modes form an appealing basis for quantum information science [86]. We use as temporal modes in this thesis Hermite-Gaussian (HG) modes as they approximately are the eigenmodes of a PDC process [70]. These modes are field-orthogonal modes in the time-frequency degree of freedom of photons.

In Figure 2.24 the working principle of a QPG is shown. A multi-mode input pulse and a single-mode pump pulse are sent to the device. The device then selects the temporal mode of the signal that coincides with the temporal mode of the pump and up-converts it. The other temporal modes are simply transmitted. As mentioned, the QPG is built on a dispersion engineered SFG process. The wavelengths of the process are typically ~ 1540 nm for the input light and ~ 860 nm for the pump light resulting in a output wavelength of ~ 560 nm. These wavelengths are commonly used since the telecom range provides low-loss

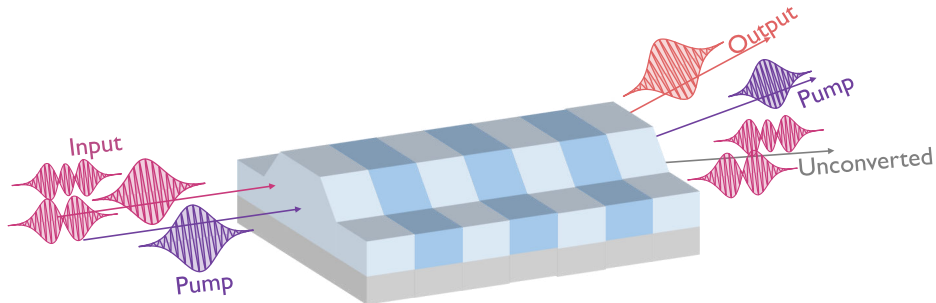


Fig. 2.24: Schematic of a QPG. The device selectively up-converts a specific mode from a multi-mode input mode with a strong pump pulse.

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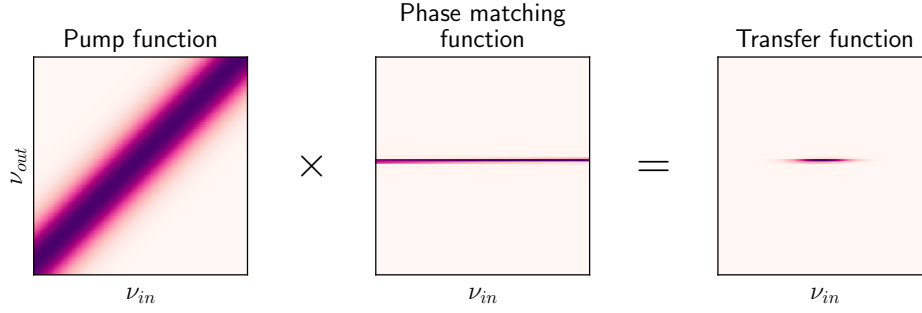


Fig. 2.25: Transfer function of a QPG. The first graph shows the pump function shaped as a HG_0 mode. The horizontal phase matching function is shown in the second graph. Multiplying both leads to the transfer function.

transmission using the existing fiber network and the pump light can easily be generated using a titanium-sapphire laser. As discussed above for an efficient SFG process, we have to fulfill energy as well as momentum conservation.

However, for QPGs these are not the only conditions that need to be fulfilled. We also need to achieve group velocity matching between the input and the pump light. This means that the input and pump light travel with the same velocity through the crystal. The reason for group velocity matching between the input and pump light becomes clear when looking at the operation principle of a QPG.

For this, we define the so called transfer function $G(\nu_{in}, \nu_{out})$ in frequency space. The function depends on the input ν_{in} and output frequencies ν_{out} and describes how the input frequencies are mapped to the output frequencies. It is given by the product of the pump function $\alpha(\nu_{in}, \nu_{out})$ and the phase matching function $\Phi(\nu_{in}, \nu_{out})$:

$$G(\nu_{in}, \nu_{out}) = \alpha(\nu_{in}, \nu_{out}) \cdot \Phi(\nu_{in}, \nu_{out}). \quad (2.142)$$

Figure 2.25 shows the transfer function for the case of a pump pulse with a Gaussian spectrum in the (ν_{in}, ν_{out}) -plane. We can see a horizontal phase matching function and the pump function in a 45° angle due to the energy conservation ($\nu_{out} = \nu_{in} + \nu_{pump}$). The transfer function is the area where both overlap. The shown phase matching function satisfies the group velocity matching condition which results in a horizontal phase matching function. The horizontal and narrow phase matching function leads to the fact that all input frequencies are converted to the same output frequency. This means that the phase matching function is independent of the input frequencies ($\Phi(\nu_{in}, \nu_{out}) \sim \Phi(\nu_{out})$). Furthermore, if the bandwidth of the phase matching function is much narrower than the input and pump pulse bandwidth, the transfer function is separable into the outer product of an input and

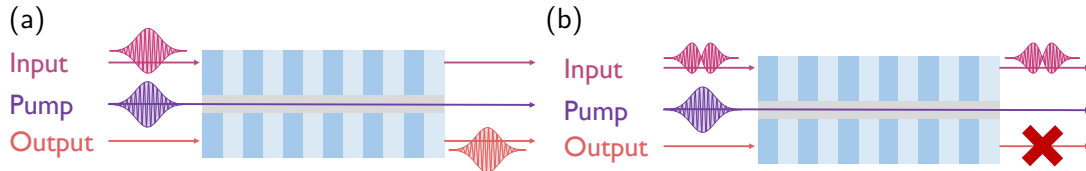


Fig. 2.26: Schematic of the working principle of a QPG. (a) If the pump and the input mode have the same mode shape, light is converted. (b) If the pump and the input mode are in different mode shapes, no conversion takes places.

output mode [87]. This means, the QPG acts as a temporal mode selective device. In other words, if we operate the QPG with a pump pulse and input pulse both shaped as HG_0 mode, the input mode will be converted (see Figure 2.26 (a)). However, if the pump pulse is a HG_0 mode and the input mode a HG_1 mode, the input mode is not converted (see Figure 2.26 (b)). The group velocity matching conditions can be determined by calculating the phase matching angle. For a SFG process the phase matching angle is given by [72]:

$$\phi_{PM} = \arctan \left(\frac{v_{g,p}^{-1} - v_{g,in}^{-1}}{v_{g,p}^{-1} - v_{g,out}^{-1}} \right). \quad (2.143)$$

For group velocity matching between the input and pump pulse, the phase matching angle is zero leading to a horizontal phase matching function. To realize a QPG in weakly guiding waveguides such as titanium indiffused waveguides in lithium niobate, one typically must utilize the host materials birefringence. Due to the low-index contrast waveguides, the dispersion cannot be engineered easily by tuning the waveguide geometry such that group velocity matching between the input and pump wavelength can only be achieved for specific wavelengths and nonlinear process types. In other words, we are limited to using type 2 processes which generally suffer from a poor optical nonlinearity. TFLN now offers the potential to explore different nonlinear processes and wavelength combinations due to its dispersion engineering capabilities which paves the way for the realization of more efficient devices, as we will later see in chapter 6.

The performance of the QPG can be assessed with two benchmarks, the temporal mode selectivity and the (normalized) conversion efficiency. The temporal mode selectivity is defined by [88, 89]:

$$S = \frac{|\rho_m|}{\sum_{n=0}^{\infty} |\rho_n|}. \quad (2.144)$$

It measures how well the device can distinguish a specific temporal mode m from all other temporal modes. Here, $\sqrt{\rho_n}$ describes the Schmidt coefficient of each mode which are normalized to $\sum_n \rho_n = 1$.

The (normalized) conversion efficiency measures how efficient the conversion is. It can be calculated as the ratio between the output light P_{out} and the input light P_{in} . By solving the coupled mode equations, see Equation 2.111, it can be shown that the efficiency of a SFG process can be estimated by [51, 89]:

$$\eta = \frac{P_{out}}{P_{in}} = \sin \left(\sqrt{\eta_{norm} P} \cdot L \right)^2. \quad (2.145)$$

Here, η_{norm} describes the normalized conversion efficiency, P the pump power and L the length of the poled area. Since the coupled mode equations in Equation 2.111 depend on the mode overlap θ of the interacting fields as well as on the nonlinear tensor element, also the normalized conversion efficiency depends on both in the following way [90]:

$$\eta_{norm} = \frac{8\pi^2 d^2 |\theta|^2}{n_{in} n_{pump} n_{out} c \epsilon_0 \lambda_{in} \lambda_{out}} \quad (2.146)$$

Here, n and λ are the refractive index and wavelength of the input, pump and output field. Thus, to achieve a higher normalized conversion efficiency, we either need to increase the mode overlap or the nonlinear tensor element. Due to the fact that the input and pump

Chapter 2. Foundations and Theoretical Concepts

pulse travel at the same speed through the waveguide, generally a high conversion efficiency can be reached in a QPG.

Now, we discussed all the necessary linear, nonlinear and quantum optical phenomena which will be used and analyzed during the scope of this thesis.

Summary

In this chapter, we introduced the theoretical concepts necessary for this thesis. First, we review the basic functionality and properties of waveguides. Afterwards, we discussed integrated devices such as directional couplers and resonators. Next, the basic principles of nonlinear optics were introduced. Moreover, we described the material platform thin-film lithium niobate. Lastly, parametric down-conversion, the Hong-Ou-Mandel effect as well as the quantum pulse gate were explained as quantum optical phenomena which will be implement in thin-film lithium niobate in the following chapters.

3

Propagation Losses in TFLN Waveguides

Waveguides are the most fundamental building block of integrated optics. For quantum applications, low-loss waveguides are crucial. When dealing with single photons every photon is important since the no cloning theorem fundamentally forbids the cloning and therefore the multiplication of quantum information carried on one photon [91]. Thus, we cannot multiply photons to preserve quantum information carried by them. Therefore, it is essential to fabricate low-loss waveguides. This creates a necessity for precise loss measurement to check if the waveguides are feasible for quantum experiments. Thus, in this chapter, we will discuss two common methods to measure propagation losses in waveguides. The first one is the Fabry-Pérot method as introduced in section 2.2.2 which is a method commonly used for standard titanium-indiffused lithium niobate waveguides in which we assume that our sample is a Fabry-Pérot resonator [44]. This method has the advantage of providing information about single waveguides as for example waveguides of single pass photon pair sources. However, it strongly depends on external parameters such as the end face reflectivity of each waveguide. For titanium indiffused lithium niobate waveguides the reflectivity can simply be calculated via Fresnel law, but for high index contrast waveguides as common in TFLN, the reflection coefficients differ from the Fresnel reflection due to the tight light confinement which will be explained later. Therefore, it would be beneficial to use a method that does not depend on any external parameter. Hence, the second method is using ring or racetrack resonators and estimating the Q-factor, as discussed in section 2.2.2. This is nowadays the common method for measuring losses of TFLN waveguides. The ring or racetrack resonator method does not provide information about each waveguide on the sample but about the combination of several waveguides on the sample. However, the method is independent from external parameters. Thus, both methods have their advantages and disadvantages and here we want to compare both to investigate the difference between them.

Therefore, this chapter is all about propagation loss measurements in TFLN waveguides. For this, we will first discuss the state of the art as well as the history of propagation losses in TFLN waveguides. Afterwards, we will introduce the setup used for all loss measurements. Next, we will briefly present the samples used to compare both methods. Finally, we will show and compare the measurement results.

Chapter 3. Propagation Losses in TFLN Waveguides

3.1 State of the Art - Propagation Losses

The first paper mentioning propagation loss measurements of TFLN waveguides was published by Hu et al. in 2009 [15]. They measured losses using the Fabry-Pérot method of 6.3 dB/cm (TE) and 7.5 dB/cm (TM). The ridge waveguides had a thin-film thickness of 730 nm, an etching depth of 460 nm and a top width of 2 μm fabricated by dry etching. Diced waveguides with a thin-film thickness of 1000 nm and a top width of 700 nm were characterized by Volk et al. in 2016 [92]. Utilizing the same measurement method, they achieved losses of 1.2 dB/cm (TE) and 2.8 dB/cm (TM) reducing the losses by a factor of 3. This reduction could be due to the different fabrication process where at this time dicing lead to smoother sidewalls. However, diced waveguides are limited in the shape and for example directional couplers or ring resonators are difficult to achieve with dicing. One year later, Wang et al. demonstrated dry etched rib waveguides with losses of 3 dB/cm using a 400 nm thin-film and a top width of 600 nm halving the losses of Hu et al [93]. But, for quantum applications these losses were still relatively high and thus, TFLN was at that moment not feasible for integrated quantum optics.

A milestone for low-loss waveguides in TFLN were the results of Zhang et al. in 2017 [18]. They demonstrated that it is possible to fabricate low-loss TFLN waveguides via dry etching. For the waveguides, they etched 350 nm from a 600 nm thin-film thickness with a top width of 2.4 μm . Using racetrack resonators, they measured losses of only 0.027 dB/cm which was two orders of magnitude lower compared to other published values of waveguide loss. This made TFLN a highly promising material for nonlinear and electro-optic applications, but also for integrated quantum optics. In the following years, many other groups published measured waveguide losses in the same order of magnitude [94, 95, 96]. Important to name is also the article by Krasnokutskaya et al. which demonstrates losses of 0.4 dB/cm in a single-mode waveguide with a top width of 810 nm, fabricated using dry etching, and losses measured with the Fabry-Pérot method [97].

Four years later, in 2021, another step towards ultra low-loss waveguides in TFLN was made when Gao et al. demonstrated waveguides with losses of only 0.0028 dB/cm via ring resonators measurements [98]. They fabricated the waveguides using femtosecond laser writing and chemical-mechanical polishing. The waveguides had a height of 720 nm and a top width of 2.11 μm . Two years later, Li et al. demonstrated losses of 0.01 dB/cm, measured via ring resonators, on a thinner lithium niobate layer of 645 nm with a top width of 1.98 μm also fabricated via femtosecond laser writing and chemical-mechanical polishing [99]. These are impressive results, however the fabrication method has drawbacks compared to dry etching. Using femtosecond laser writing limits the resolution and thus nanostructured devices are challenging to fabricate. Moreover, achieving good reproducibility is demanding. Thus, using dry etching, as a waveguide fabrication method, especially when thinking about wafer-scale production, is most likely the better choice.

Hence, it was a big step forward for integrated quantum optics in TFLN, when, in 2024, Zhu et al. realized waveguides with losses of only 0.013 dB/cm using a 600 nm thin-film, an etching depth of 325 nm and a top width of 3 μm using dry etching [100]. These losses are only one order of magnitude higher than the intrinsic losses of the material [35]. Thus, ultra low-loss waveguides feasible for integrated quantum optics experiments can be realized in TFLN.

To get an overview of the evolution of the loss values, measurement method and the fabrication technique, we plot the losses as a function of the publication year in Figure 3.1.

3.1. State of the Art - Propagation Losses

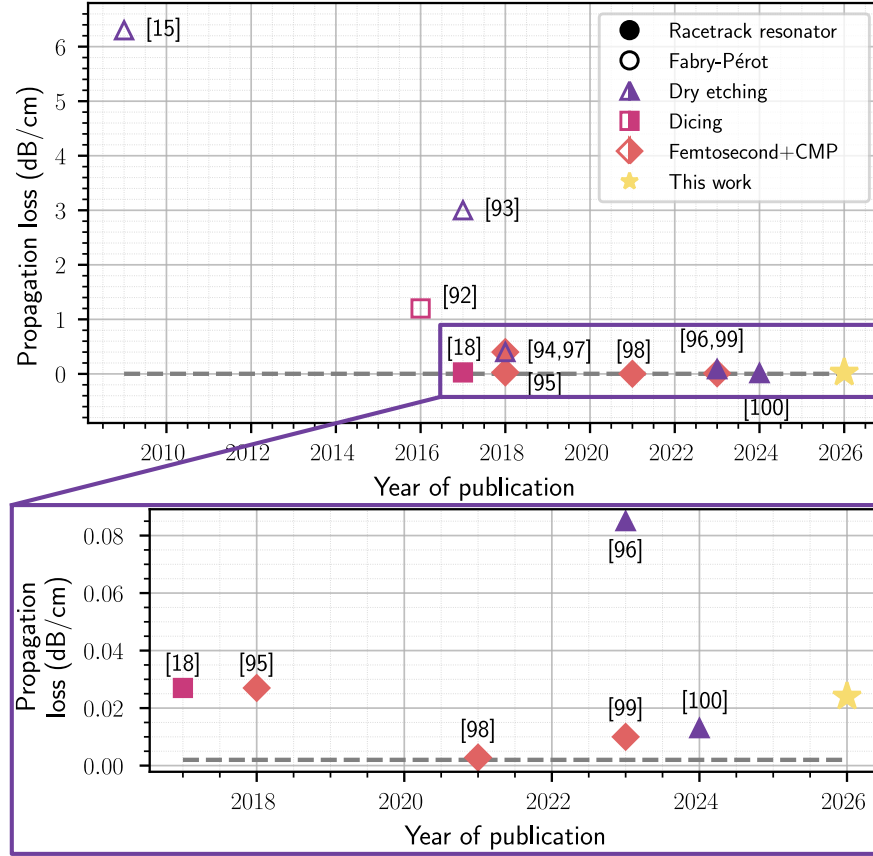


Fig. 3.1: Propagation loss as a function of the publication year. If the marker is filled, the losses are measured with the racetrack method. The Fabry-Pérot method is used, if the marker is not filled. For each fabrication method, we use a different color and marker. The graph shows that the losses decreased over the years and that low-loss waveguides can be fabricated with every fabrication method.

The not filled markers represent losses measured with the Fabry-Pérot method and the filled markers are for the ring or racetrack method. For each fabrication method, we use a different colored marker. Moreover, we indicate the material limit by a grey dashed line. We can see that by improving the fabrication process, the propagation losses of TFLN waveguides decreased from around 6 dB/cm to only 0.01 dB/cm. Moreover, it shows that with each fabrication method, we can produce low-loss waveguides. This indicates that TFLN is an ideal material platform for ultra-low loss waveguides and due to the high second-order nonlinearity and the electro-optical coefficient also a promising material platform for integrated quantum optics. However, we can also see that the loss measurement method changed from the Fabry-Pérot method to using ring resonators and that the investigated waveguides have a rather broad top width of more than $2\text{ }\mu\text{m}$. The ring resonator method has the drawback that this only gives losses of ring structures and not of single straight waveguides. But, for the Fabry-Pérot method the reflectivity of the waveguide is necessary. Thus, in this chapter we want to compare the two methods and analyze the difference between the two methods as well as compare our losses to the state of the art losses. Furthermore, we want to investigate the influence of the geometry namely the top width on

Chapter 3. Propagation Losses in TFLN Waveguides

the losses since most of the waveguides used in publications have a rather broad top width. For this, we will first start with describing the setup for these measurements.

3.2 Loss Characterization Setup

For analyzing the propagation losses of TFLN waveguides with the Fabry-Pérot and the racetrack resonator method, we use the same setup which is shown in Figure 3.2. The setup

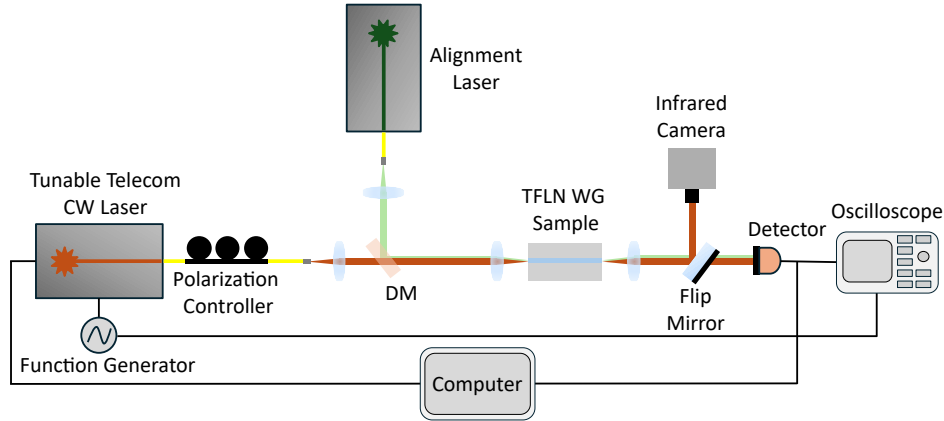


Fig. 3.2: Schematic of the setup to measure propagation losses of TFLN waveguides.

consists of two laser sources. The first one is an alignment laser with a visible wavelength, at either green or red. The second laser source is a CW-laser in the telecom range. The wavelength-tunable CW-laser is connected to a fiber polarization controller using a SMF-28 fiber to set the polarization. Both lasers are coupled to free-space using a lens. The two laser beams are combined with a dichroic mirror (DM) and propagate collinearly afterwards. The light is coupled into and out of the TFLN waveguide using aspheric lenses. After the sample, the light is measured with a power meter or photodiode. Both are connected to a computer as well as the laser. For the measurement of the racetrack resonators also an oscilloscope and a signal generator are used which are connected to the photodiode and to the laser. This is done to scan the wavelength in a finer manner. For easy first alignment an infrared camera can be used by putting a mirror after the sample to image the output of the waveguide as well as to characterize directional couplers. This setup is used to measure propagation losses with the Fabry-Pérot method and the racetrack resonator method.

3.3 Loss Measurements

In this part, we will first show the sample layout we used to measure and compare the losses with both methods. Afterwards, we will discuss the results of the Fabry-Pérot method and of the racetrack resonator method. Finally, we will compare both methods and investigate the influence of the geometry on the losses. All measurements were performed with TE polarized light.

3.3. Loss Measurements

3.3.1 Sample Layout

We want to measure and compare the propagation losses using the Fabry-Pérot and racetrack resonator method. Moreover, we also want to investigate the influence of the geometry on the losses. For this purpose, we design three samples with three different top widths. The geometry of the waveguide is a 600 nm thin-film thickness, a etching depth of 300 nm, a sidewall angle of 30° and a SiO_2 cladding layer with a $1\text{ }\mu\text{m}$ thickness. The top width is on the first sample $1\text{ }\mu\text{m}$, on the second sample $2\text{ }\mu\text{m}$ and on the third sample $3\text{ }\mu\text{m}$. With these three top widths, we can measure the losses of a single mode configuration ($1\text{ }\mu\text{m}$ top width) as well as a multi mode configuration ($3\text{ }\mu\text{m}$ top width) to be able to analyze the influence of the geometry on the losses. In Figure 3.3 a laser microscope image of the sample with a $3\text{ }\mu\text{m}$ top width is shown. In the middle of the sample we have a big waveguide with a width of $20\text{ }\mu\text{m}$ (marked in pink) which is used for alignment as well as two directional coupler to check the coupling behavior. Below and above the alignment waveguide there are on each side six groups of waveguides. Each group consists of a straight waveguide and a racetrack resonator. Thus, overall we have twelve waveguide groups on the sample with six different racetrack length such that each length is twice on the sample.

The bending radius of the racetrack resonators is $200\text{ }\mu\text{m}$ and we use a circular bending. On the sample with a $3\text{ }\mu\text{m}$ top width a Euler bending is used to suppress coupling to higher order modes in the bending [101]. The directional couplers are used to measure the coupling behavior if wanted. The coupling length is chosen in such a way that the racetrack resonator is undercoupled, as discussed in section 2.2.2. Designing the racetrack in this way makes the loss evaluation easier since it is hard to determine if the resonator is under- or overcoupled if the racetrack resonator is closed to critical coupling. Thus, trying to achieve

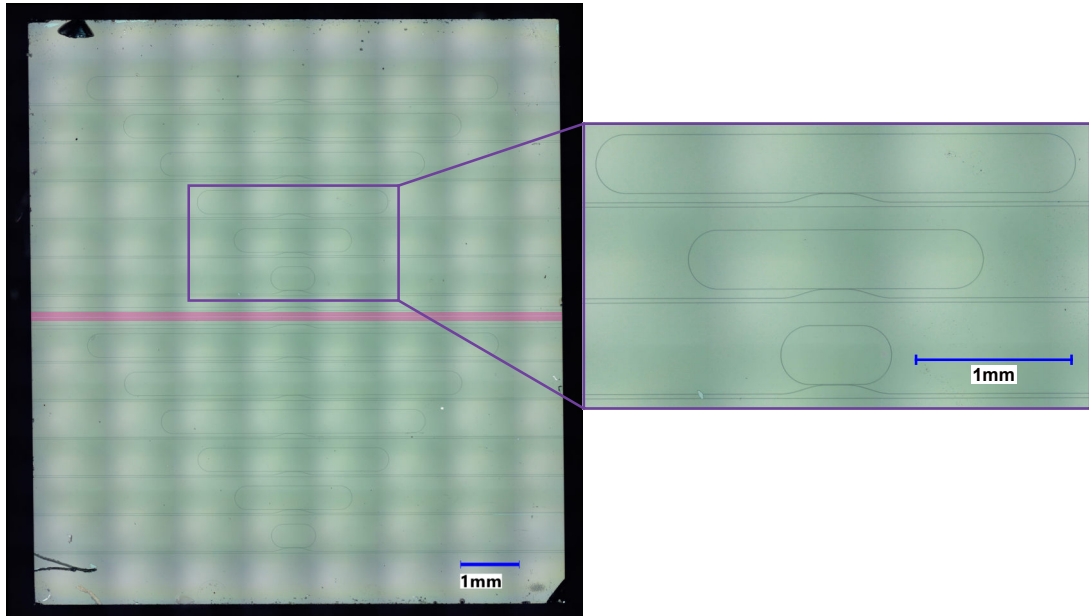


Fig. 3.3: Overview of a sample used to measure the losses with the racetrack method and Fabry-Pérot method with a $3\text{ }\mu\text{m}$ top width. We mark the alignment waveguide in pink. The enlarged image shows three racetracks with different sizes and next to each racetrack a straight waveguide.

Chapter 3. Propagation Losses in TFLN Waveguides

undercoupling helps to prevent to be at the point of critical coupling. Based on this, the racetrack resonators have the following parameters:

- 1 μm sample:
 - center-to-center gap: 2.3 μm
 - coupling length: 100 μm
 - straight length of racetrack: 370 μm , 1616 μm , 2862 μm , 4108 μm , 5354 μm , 6600 μm
- 2 μm sample:
 - center-to-center gap: 3.3 μm
 - coupling length: 150 μm
 - straight length of racetrack: 350 μm , 1600 μm , 2850 μm , 4100 μm , 5350 μm , 6600 μm
- 3 μm sample:
 - center-to-center gap: 4.3 μm
 - coupling length: 600 μm
 - straight length of racetrack: 900 μm , 2040 μm , 3180 μm , 4320 μm , 5460 μm , 6600 μm .

With these samples we can measure the losses of the straight waveguides using the Fabry-Pérot method and on the same sample we can use the racetrack resonators to also estimate the losses with this method. Therefore, we are able to compare both methods on the same sample. Moreover, we can also compare the losses of the three different geometries from the three different samples. Next, we will discuss the loss measurements with the Fabry-Pérot method.

3.3.2 Fabry-Pérot Measurements

To measure the losses with the Fabry-Pérot method, we use the setup as shown in Figure 3.2. We scan the CW-laser from 1550 nm to 1550.7 nm in 2 pm steps and measure the transmitted power with a power meter connected to the computer. To estimate the losses, we calculate the contrast K from the smoothed data. As an end face reflectivity R we use the standard Fresnel reflectivity of lithium niobate of 14.2 %. Together with the sample length L we can estimate the losses.

In Figure 3.4 we show the transmitted power as a function of the wavelength for one waveguide of each sample. We can see uniform oscillations. Even for the broader top widths, no beating between different modes is visible in the oscillations meaning we only excited the fundamental mode.

We measure the losses of all straight waveguides on each sample and calculate the average losses as well as the standard deviation. The average losses and standard deviations are 1 μm : (1.99 ± 1.36) dB/cm, 2 μm : (-0.31 ± 1.53) dB/cm, 3 μm (-1.10 ± 0.06) dB/cm. The large errors of the sample with 1 μm and 2 μm show a large variation of waveguide quality over the sample. This could be due to higher losses of different waveguides but also due to

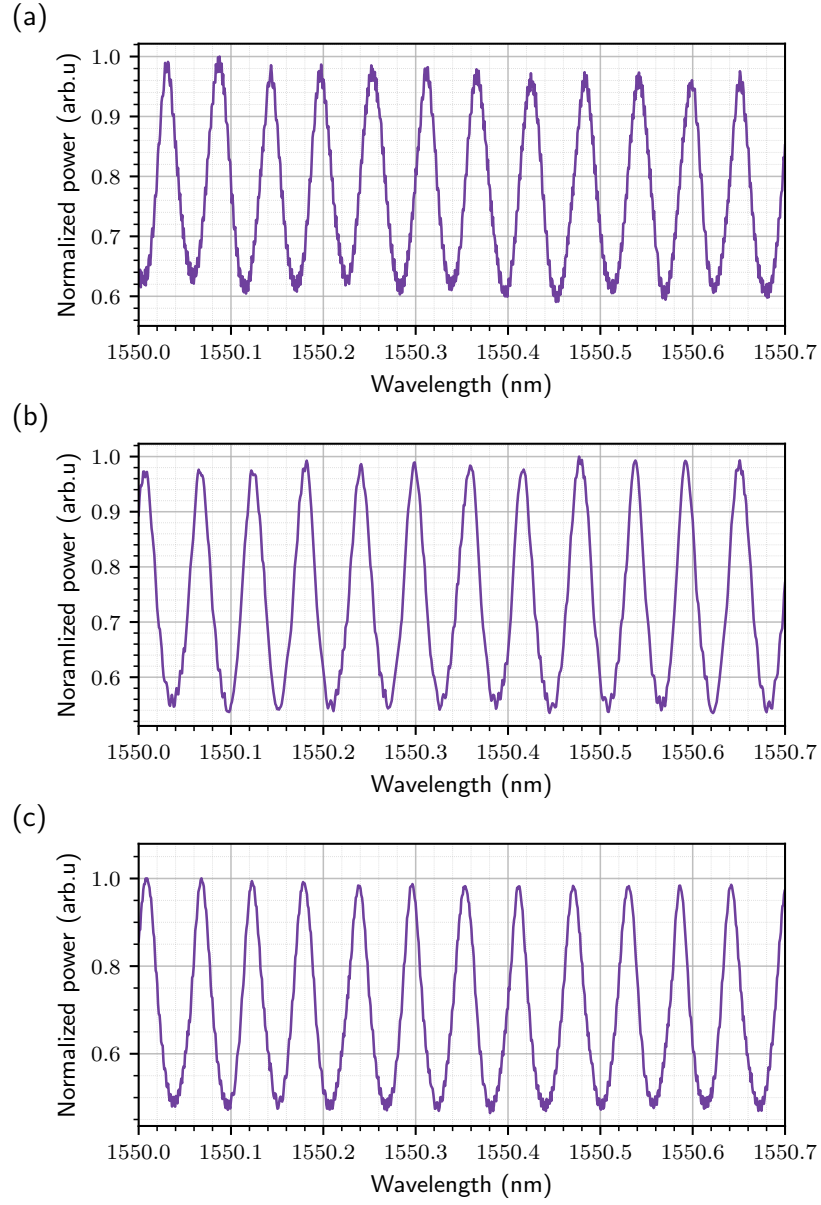


Fig. 3.4: Fabry-Pérot oscillations from waveguides with three different top widths. (a) 1 μm top width. (b) 2 μm top width. (c) 3 μm top width.

Chapter 3. Propagation Losses in TFLN Waveguides

a broken end face. This shows that it is important to check each waveguide individually. The most important aspect is that negative losses are measured for a $2\text{ }\mu\text{m}$ and $3\text{ }\mu\text{m}$ top width. Negative losses means that we would measure more light after the waveguide than before the waveguide which is not possible in single straight waveguides. The negative losses are most likely due to the end face reflectivity assuming Fresnel reflectivity. This is only valid for plane waves, but due to the small size of the waveguides which is in the range of the wavelength this approximation is not true anymore [102]. Thus, to use this method the end face reflectivity needs to be simulated for each geometry for precise loss estimation. Moreover, the losses decrease with increasing top width. This is due to the fact that the fundamental mode is more confined in the middle of the waveguide for broader top widths and thus the evanescent field has lower overlap with the rough sidewalls which is due to the etching process which causes scattering losses. However, we want to compare the Fabry-Pérot losses with the losses estimated with the racetrack resonators to analyze if we have a constant offset which we can use as a calibration factor.

3.3.3 Racetrack Resonator Measurements

To estimate the losses with the racetrack resonator method, we measure the transmitted power as a function of the wavelength for each resonator on the sample. To do this, we use the setup as shown in Figure 3.2. Here, we scan the CW-laser by applying a triangular voltage to scan the laser cavity. The transmitted light is measured with a photodiode connected to an oscilloscope.

From the measured spectra, we determine all necessary parameters as discussed in section 2.2.2 to estimate the losses. In Figure 3.5, the transmission spectra for three different racetracks with different width are shown. In all plots we can see the Fabry-Pérot oscillations due to the end faces of the sample as well as the dips from the racetrack resonator. Our system consists of two coupled resonators: a racetrack resonator that produces sharp, high-finesse, typically Lorentzian, resonance dips, and a Fabry-Pérot resonator formed by reflections at the end facets of the bus waveguide, giving rise to low-finesse oscillations in the transmission spectrum. Interaction between the two cavities can alternate the Lorentzian shape of the racetrack dips, resulting in a modified asymmetric line shape known as Fano resonances [103]. Their exact shape depends on the intra-cavity coupling strength and the relative phase between the resonant pathways [104]. To accurately extract the dips' line widths and depths of the racetrack resonances, each dip we fit with a Fano line shape, as illustrated for a representative resonance in Figure 3.6. For the $3\text{ }\mu\text{m}$ waveguide we can see a beating of different modes which we did not measure for the straight waveguides. The excitation of higher modes might be due to the bending or due to alignment. We multiply the total losses of each racetrack by its corresponding length and plot these losses as a function of the straight part of the racetrack (see Figure 3.7). Then, we estimate the losses by applying a linear fit for each waveguide width. Here, we can see the same trend as for the Fabry-Pérot method which is that the losses decrease with increasing top width. For $1\text{ }\mu\text{m}$ we estimate losses of $(1.462 \pm 0.195)\text{ dB/cm}$, for $2\text{ }\mu\text{m}$ $(0.070 \pm 0.016)\text{ dB/cm}$ and for $3\text{ }\mu\text{m}$ $(0.024 \pm 0.036)\text{ dB/cm}$. It is noteworthy, that the error for the $1\text{ }\mu\text{m}$ sample is relatively small compared to the $2\text{ }\mu\text{m}$ and $3\text{ }\mu\text{m}$ sample where the error is the same order of magnitude as the losses. This is due to the fact that the losses of the broader waveguides are so low that the length difference does not lead to significant change in the total losses (see Figure 3.7). Thus, small deviations, e.g., from fabrication have a considerable influence on the losses and the fitting of the data. Hence, to have precise knowledge of the losses for these top width a

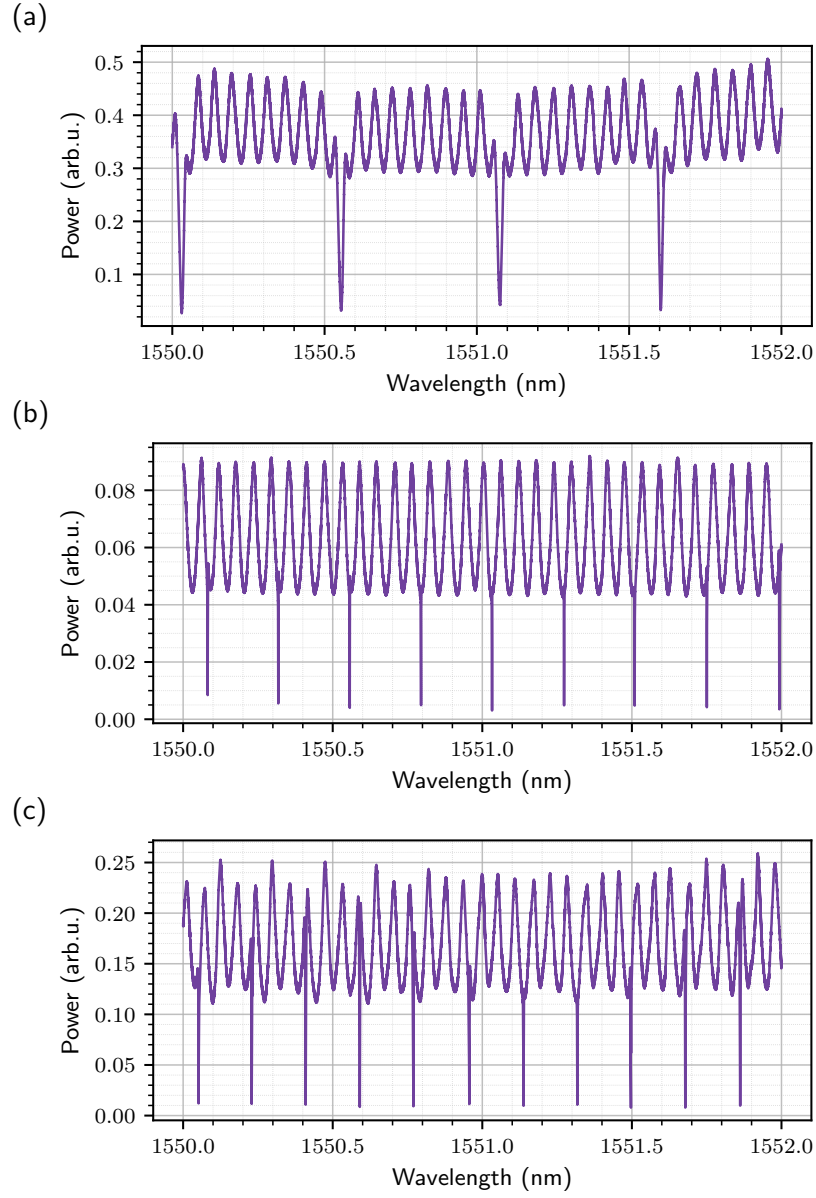


Fig. 3.5: Transmission spectrum of racetrack resonator with different top widths. (a) $1\ \mu\text{m}$ top width. (b) $2\ \mu\text{m}$ top width. (c) $3\ \mu\text{m}$ top width.

Chapter 3. Propagation Losses in TFLN Waveguides

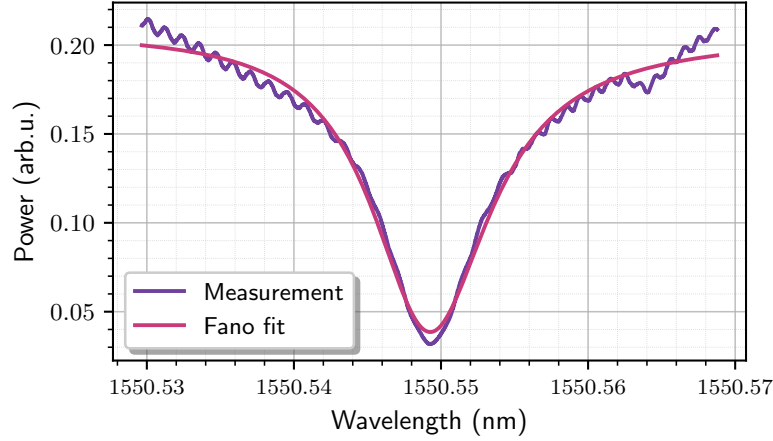


Fig. 3.6: Power as a function of the wavelength for one resonance dip of a racetrack resonator. We fit a Fano resonance to the dip to retrieve the necessary parameters.

more pronounced length difference between the racetracks is necessary.

3.3.4 Comparison

As the last part of this chapter, we want to compare the two characterization methods as well as our best measured losses to the state of the art. For this, we plot the losses measured with both methods as a function of the top width (see Figure 3.8). For better visibility, we connected the measured losses for each method.

At first glance, we can see that both methods show the same trend of decreasing losses with increasing top width. As discussed before this is due to the fact that the mode is more confined in the middle of the waveguide and the evanescent field has less overlap with the sidewalls. This reduces scattering losses at the rough sides of the waveguide from the etching process. We also validated this behavior theoretically with the help of Dr. Manfred Hammer [105]. This also indicates that for low-loss routing on chip broader waveguides should be used. Only in interaction regions such as poled areas where single mode interaction is crucial to avoid higher order processes the waveguide top width should be reduced.

For small top widths and thus higher losses both methods yield similar results. However, if the losses decrease the Fabry-Pérot method provides negative losses. For a $2\text{ }\mu\text{m}$ top width the Fabry-Pérot losses are already negative, but still somewhat close the racetrack resonator losses. If the losses decrease further, the results from the two method drift further apart and the Fabry-Pérot losses are below -1 dB/cm . This means the Fabry-Pérot measurement is suitable for a rough estimation of high losses, but not suited for low-loss measurement of TFLN waveguides. Moreover, the reflectivity also depends on the quality of the end faces of the waveguide. A broken end face has a different reflectivity compared to a intact one which can lead to inaccurate measured losses. Moreover, the end face cannot only be broken but it can also be tilted which reduces the contrast and thus influences the measured losses as well. Thus, even when having precise knowledge about the assumed reflectivity, the losses are still prone to errors of the real reflectivity of a sample and the quality of the end faces. The racetrack resonator method is more precise especially when dealing with ultra low loss waveguides but for this a more pronounced length different in the racetracks is necessary. The varying offset between the two methods also shows that the reflectivity is different for

3.3. Loss Measurements

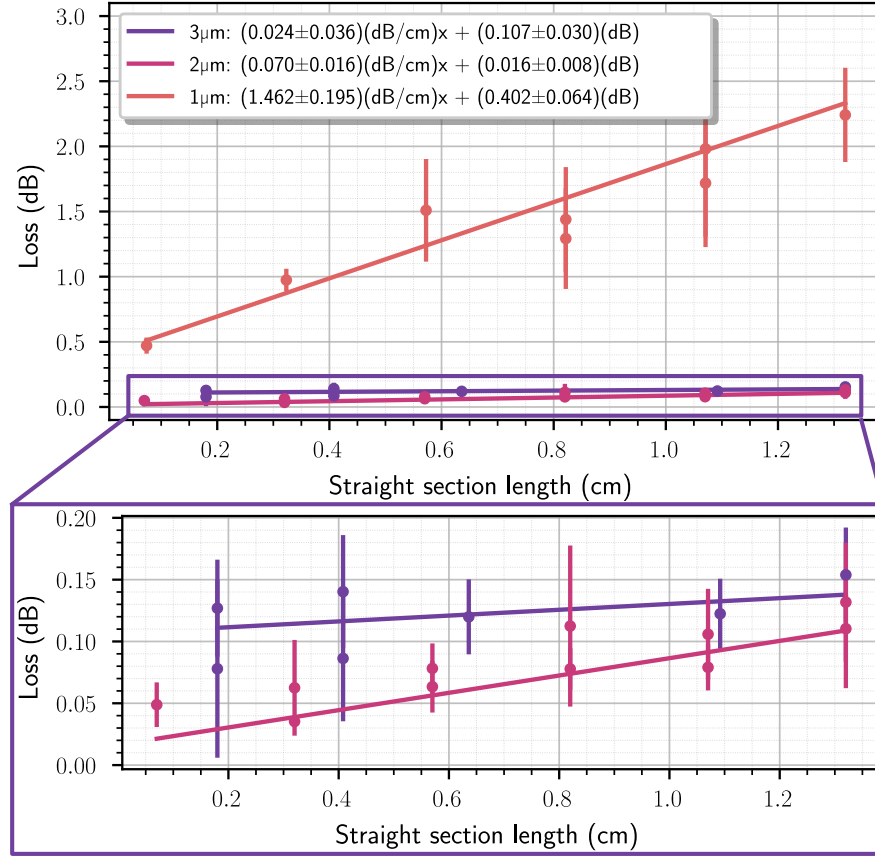


Fig. 3.7: Total losses as a function of the straight waveguide section for all three top widths. For each top width, we applied a linear fit to determine the propagation losses from the slope and the bending losses from the y-intersection. For better visualization the lower plot shows an enlarged view of the lower part of the upper plot.

each geometry. This varying reflectivity is due to the small size of the waveguide which is in the sub-wavelength regime. Therefore, we cannot assume plane waves anymore and Fresnel law is not applicable. This leads to the fact that the reflectivity is slightly different for each geometry. One approach is to change the reflectivity such that the Fabry-Pérot losses and the racetrack resonator losses are the same. To match the losses of the 1 μm sample measured with the Fabry-Pérot method to the racetrack losses a reflectivity of around 13 % is needed. However, due to the large variation of losses on the sample a reflectivity of 8.4 % is required for the worst waveguide and 16.5 % for the best waveguide. For the 2 μm sample the reflectivity changes to 15.2 %, but can vary between 7.4 % and 23.3 % to match the losses of the two methods. A reflectivity of 17.5 % is needed for the 3 μm sample. Only for this sample the variation in the reflectivity is in the range of 0.4 %. Therefore, only for the 3 μm we can estimate the necessary reflectivity. For the other two top width more samples are needed to estimate the reflectivity. But, we can also observe that the reflectivity is below 14.2 % for the 1 μm sample and above for the other two samples. This could be due to the fact that the mode "sees" more SiO₂ for the smallest waveguide compared to the broader ones. The mode of a waveguide with a 1 μm top width leaks more into the SiO₂ layer beneath and above the waveguide than the mode of a 3 μm top width waveguide for

Chapter 3. Propagation Losses in TFLN Waveguides

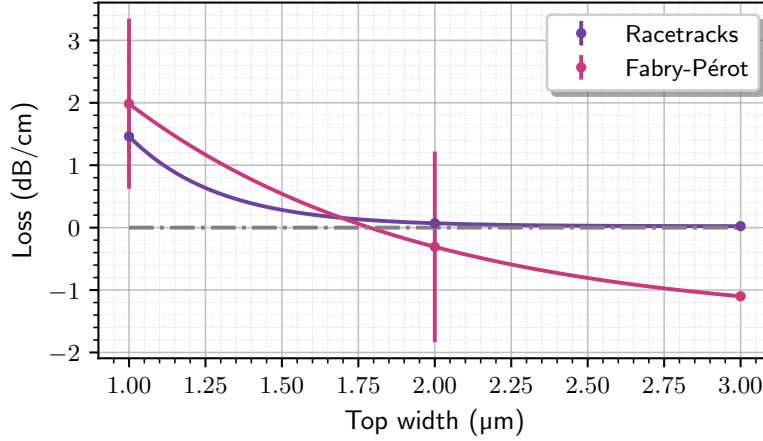


Fig. 3.8: Comparison of the losses measured with the Fabry-Pérot method and the racetrack resonator method as a function of the top width. The losses decrease with increasing top width as expected. The losses measured with Fabry-Pérot method get negative for 2 μm and 3 μm top width. We can see an decay of the losses measured with the racetrack resonators. We connected the measurement points for better visibility and added a grey dashed line at 0 dB/cm.

which the mode is more confined in the center of the waveguide. Thus, the mode of the smaller waveguide is more influenced by the SiO₂ layers. Since SiO₂ has a lower refractive index than lithium niobate, the reflectivity for a smaller waveguide might be lower than for a broader waveguide. Thus, the reflectivity is still a parameter under investigation. As a next step, the reflectivity could be investigated theoretically via simulations to compare it to the measurements. Moreover, more samples should be characterized to have more statistics. Lastly, we can see a significant decrease in the losses from single-mode to multi-mode waveguide structures. This also means when comparing the losses to the literature one has to closely compare the waveguide geometry. Moreover, this also should be considered when designing devices. Comparing our best measured losses to the state of the art, our losses are in the same order of magnitude (see Figure 3.1). This means, we are able to fabricate and characterize low-loss waveguides in TFLN.

In conclusion, the Fabry-Pérot method is not practical to measure ultra-low loss waveguides with the standard reflectivity. To use this method, a more in-depth analyses of the geometry dependent reflectivity is necessary. The racetrack resonator method is a feasible method, but it cannot measure the losses of single straight waveguides. Thus, it would be beneficial to match the Fabry-Pérot and racetrack losses by a detailed analysis of the end face reflectivity to be able to analyze single straight waveguides precisely. However, we successfully implemented the ring resonator method and showed that we are able to fabricate ultra low-loss multi-mode waveguides with losses of only (2.4 ± 3.6) dB/m which are on par with state of the art losses.

Summary

In this chapter, we discussed how propagation losses in TFLN waveguides can be estimated. For this, we compared two common methods. The first method is the Fabry-Pérot method which is the standard method to measure losses of titanium indiffused lithium

3.3. Loss Measurements

niobate waveguides which provides information about single waveguides. Secondly, we realized the racetrack resonator method which is a typical method for loss estimation of TFLN waveguides. However, with this method, only the losses of racetracks can be measured and no information about specific single waveguides can be gained. The Fabry-Pérot method is not usable for TFLN waveguides without an in-depth analysis of the reflectivity. With the racetrack resonator method we measured the expected decay in the propagation losses with increasing top width. This leads to the conclusion that broader waveguides should be used for routing the light on chip and only in interaction regions the top widths should be reduced to avoid higher order processes to achieve less losses on chip. We showed that we are able to fabricate and characterize ultra low-loss waveguides in TFLN with loss of (2.4 ± 3.6) dB/m for multi-mode waveguides which is line with state of the art losses.

4

Hong-Ou-Mandel Interference in a TFLN Directional Coupler

In the last chapter, we have seen that low-loss waveguides are feasible in TFLN making it a highly promising material platform for integrated quantum optics. Thus, we want to use it to show a genuine quantum effect. We choose the most fundamental quantum phenomenon namely the interference of two single photons at a balanced beam splitter also known as Hong-Ou-Mandel (HOM) interference. Therefore, we will discuss in this chapter the demonstration of the HOM effect in an integrated beam splitter in TFLN. We start with explaining the history and state of the art of integrated HOM interference. Next, we will show the modeling as well as the characterization of a directional coupler. Afterwards, we show the setup built to measure the HOM dip. Lastly, we discuss the measurement results. The following results have been published in [106].

4.1 State of the Art - HOM Interference

The interference at a balanced 50:50 beam splitter (as discussed in subsection 2.5.2) was first shown in 1987 by C.K. Hong, Z.Y. Ou and L. Mandel [77]. Since then it is one of the simplest but also one of the most fundamental quantum optical phenomena [107]. It is a fundamental building block of quantum networks and lies at the heart of various quantum optical experiments such as boson-sampling [108] or quantum key distribution [109]. Hong, Ou and Mandel used a bulk beam splitter, but to realize scalable quantum experiments an integrated beam splitter would be beneficial. A beam splitter can be integrated in different material platforms and thus, HOM interference has already been shown in numerous materials such as silicon-on-insulator [110, 111, 112], laser written waveguides in glass [113] or silicon nitride [114].

Most of the platforms were $\chi^{(3)}$ -materials which can have the advantage of small device footprints due to the high index contrast waveguides. However, those materials lack the possibility of generating single photons efficiently via three wave mixing processes which are essential for fully integrated quantum experiments.

Therefore, using lithium niobate would be beneficial and a fully integrated HOM experiment was also shown in titanium indiffused waveguides in lithium niobate [8]. But, the device was relatively bulky due to the large bending radius of the low index contrast waveguides. To achieve a directional coupler in a $\chi^{(2)}$ -material with a high integration density, TFLN is an optimal platform to miniaturize and enable integrated quantum circuits.

To integrate a beam splitter in TFLN or any other material platform, different approaches

HOM Interference in a TFLN Directional Coupler

can be utilized such as a directional coupler or a multi-mode interference coupler. In TFLN various implementations have been published [115, 116, 117, 118, 119, 120]. Here, we use a directional coupler scheme in TFLN to demonstrate HOM interference of two photons at telecom wavelengths.

Following our publication, numerous papers have used HOM interference in TFLN on chip for integrated quantum optical experiments such as HOM interference in a Mach-Zehnder modulator [121, 122], HOM interference between two independent PDC sources on chip [123, 124] or HOM interference in a rapid adiabatic coupler [125].

4.2 Modeling of a TFLN Directional Coupler

To show the HOM interference in TFLN, we first need to design a directional coupler, which we introduced in subsection 2.2.1, that meets the following requirements. We want to realize a directional coupler consisting of waveguides that only guide the fundamental TE mode at 1550 nm to suppress coupling to higher order modes. Moreover, a short coupling length is beneficial to reduce losses and achieve a small footprint. Lastly, the coupler should be broadband to support the bandwidth of the single photons used for the interference. Based on these requirements, we search for a suitable coupler geometry while taking the following fabrication constraints at that time into account: the thin-film thickness is fixed to 600 nm, due to the etching process, we have a sidewall angle of 60° . Moreover, to protect the sample, we apply a cladding layer of 150 nm.

We start with investigating the single-mode behavior of a single straight waveguide. For this, we use the simulation model described in section 2.4.2 and simulate for 1550 nm the guided TE modes. To find a single-mode waveguide, we first fix the top width to $1\text{ }\mu\text{m}$ and scan the etching depth. For each etching depth, we look at the number of modes simulated in the *MODE* solver of *Lumerical*. We scan the etching depth from 50 nm to 300 nm. During the simulations, we can observe that with increasing etching depth the waveguide changes from a single-mode to a multi-mode configuration. This is due to the higher confinement of the modes inside the resonator which leads to the support of higher order modes. For the TE mode, we can see that the waveguide gets multi-mode around 200 nm. Based on this, we choose a etching depth of 150 nm to include possible fabrication deviations to assure that we stay in the single-mode regime.

Next, we investigate the gap G of the directional coupler. Here, we define the gap as the distance from the middle of one waveguide to the middle of the other waveguide meaning from center to center. For this, we extend our simulation of a single straight waveguide by adding a second waveguide. This resembles the cross-section of a directional coupler and the symmetric and anti-symmetric mode can be simulated. Here, it is important that both modes, the symmetric and anti-symmetric mode, are guided in the waveguides. We scan the gap from $1.5\text{ }\mu\text{m}$ to $4\text{ }\mu\text{m}$ and calculate for each gap the crossover length L_x according to Equation 2.36. We plot the results in Figure 4.1 which shows the crossover length as a function of the gap. Here, we can observe that the crossover length increases exponentially with increasing gap. This is due to the fact that modes inside the two waveguides couple less and less with each other the larger the distance between the two waveguides. For a small center-to-center gap of $1.5\text{ }\mu\text{m}$ there is only a distance of $0.5\text{ }\mu\text{m}$ from the top edge of one waveguide to the top edge of the second waveguide. Due to the waveguide angle, the waveguides have a bottom width of $1.52\text{ }\mu\text{m}$. This leads to the fact that the two waveguides are not well separated if they are only separated by $1.5\text{ }\mu\text{m}$ from the center of the first

4.2. Modeling of a TFLN Directional Coupler

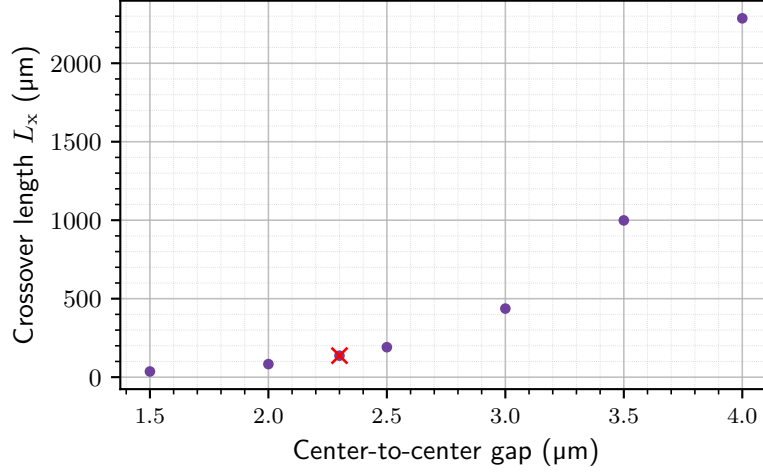


Fig. 4.1: Crossover length L_x as a function of the center-to-center gap. With increasing gap, the crossover length increases exponentially. The red cross marks the chosen coupling length for the directional coupler design.

waveguide to the center of the second waveguide. Due to this the modes merge together in the gap. Between a gap of $2\mu\text{m}$ and $2.5\mu\text{m}$ the waveguides get well separated and the modes are mostly guided in their respective waveguides. Increasing the gap even more leads to an increase of the crossover length. Since we want to realize a small footprint device to take advantage of the high confinement waveguides in TFLN, we want to use a small gap. However, due to the resolution of the laser lithography which is used to fabricate the directional couplers, we are limited to a minimal distance between the bottom edges of the waveguide of around $1\mu\text{m}$. Based on this, we decide on a gap of $2.3\mu\text{m}$. This gap leads to a coupling length of $137\mu\text{m}$ calculated via Equation 2.36 and the simulated effective refractive indices of the symmetric and anti-symmetric modes. In Figure 4.2 we plot the fundamental TE mode as well as the symmetric and anti-symmetric mode of the chosen

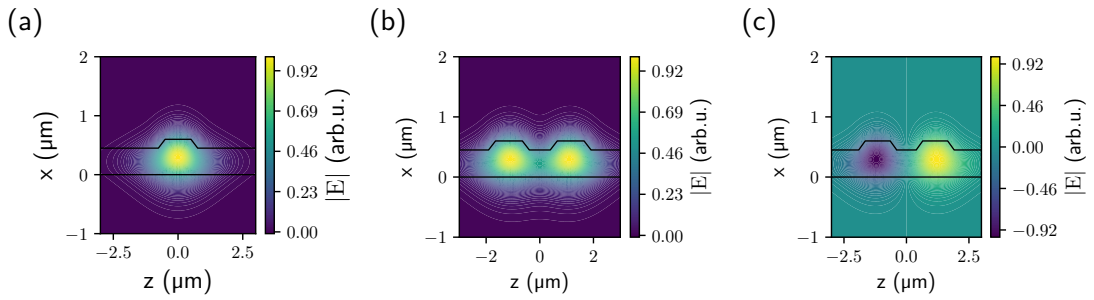


Fig. 4.2: Simulated mode distributions of the cross-section of a single waveguide and a directional coupler. Plot (a) shows the mode distribution of the cross-section of a single waveguide. In plot (b) and (c) the mode distribution of a cross-section of a directional coupler consisting of two waveguides next to each other are shown. The symmetric mode of the directional coupler is shown in (b) and the anti-symmetric in (c).

HOM Interference in a TFLN Directional Coupler

geometry ($t=600$ nm, $h=150$ nm, $w=1$ μm , $a=60^\circ$, $c=150$ nm, $G=2.3$ μm).

Lastly, we also check the bandwidth of the directional coupler at 50:50 splitting ratio. For this, we simulate the coupler for a range of wavelengths and calculate from the effective refractive index of the symmetric and anti-symmetric mode the splitting ratio for a fixed coupling length of half of the crossover length of 137 μm . In Figure 4.3 the splitting ratio is plotted as a function of the wavelength. It can be seen that in a wavelength range from 1540 nm to 1560 nm the splitting ratio changes less than 2 %. The single photons used for the experiment have a bandwidth of 1.8 nm at 1545 nm. Thus, the chosen directional coupler geometry is ideally suited for measuring HOM interference.

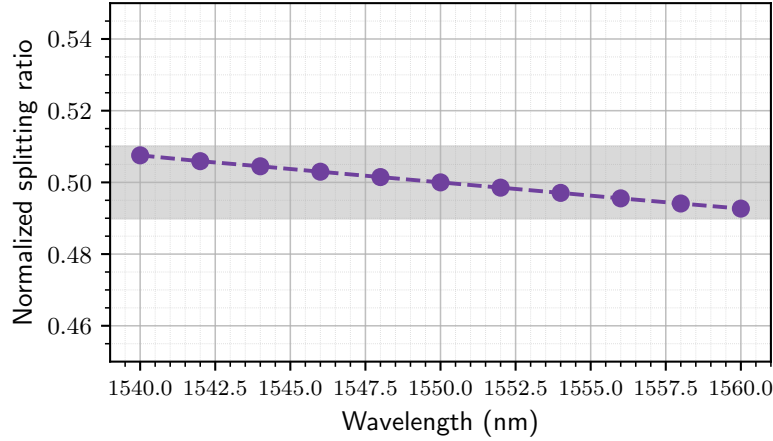


Fig. 4.3: The splitting ratio as a function of the wavelength. In grey, we mark a deviation of $\pm 1\%$ from the 50:50 splitting ratio. It can be seen that the directional coupler is broadband enough for the used photons with a bandwidth of 1.8 nm.

Based on the coupler geometry we design a sample to investigate the coupling behavior by varying the interaction length. The input and output waveguides have a separation of 60 μm at the beginning and end. To bring them close for coupling, a 300 μm cosine bending is used. We change the interaction length from 30 μm to 580 μm .

The fabricated geometry deviates slightly from the simulated one. The thin-film thickness is 605 nm, the etching depth is 130 nm and the top width 1.066 μm . All other parameters are as simulated. The fabricated geometry has a simulated crossover length of 120 μm . Next, we characterize the couplers to find the best coupler with a 50:50 splitting ratio that we can use for the HOM experiment.

4.3 Characterization of Directional Couplers

To analyze the directional coupler, we use the setup shown in Figure 3.2 by putting in the flip mirror to guide the light to the infrared camera to measure the modes of the coupler. To do so, we couple a TE polarized 1550 nm CW-laser into one of the input waveguides. After the sample, we measure the two modes of the directional coupler with an infrared camera. In Figure 4.4, we plot the measured mode pictures of two directional couplers with different lengths. The upper one shows only one mode meaning all the light stayed in the input waveguide. The lower picture shows two modes with equal intensity meaning half of the light coupled from one waveguide to the next. We take such mode pictures for all couplers for

4.3. Characterization of Directional Couplers



(a) Interaction length: 58 μm .



(b) Interaction length: 257 μm .

Fig. 4.4: Measured mode pictures of the two output ports of two different directional couplers. Depending on the length of the coupler (a) all of the light stays in the initial waveguide or (b) couples between the two waveguides.

both input waveguides. Then, we calculate the splitting ratio using the following equation:

$$\eta = \frac{P_1}{P_1 + P_2} \quad (4.1)$$

where $P_{1,2}$ is the power of the left and right mode. To get the power of each mode, we integrate over the pixels of each mode. In Figure 4.5, we plot the splitting ratio as a function of the interaction length. We can see that the light oscillates between the two waveguides. But, we can also observe that even for no interaction length light couples between the two waveguides. This offset occurs because the light already starts to couple in the bendings from one waveguide to the other. Moreover, we can observe that not all splitting ratios of waveguide I and waveguide II add up to 1. We attribute this to asymmetric losses in the two output waveguides.

To determine the crossover length, we apply a sinus fit to the measured data. From the two fits, we obtain slightly different crossover lengths for waveguide I and waveguide II of 114.8 μm and 110.7 μm . This deviation is due to differences in the waveguide configuration. We assume that the crossover length is the average of both measured lengths and the error is the standard deviation leading to a crossover length of $L_x = (112.75 \pm 2.89) \mu\text{m}$ which is close to the simulated one of 120 μm . Deviations could be due to the fact that we did not include bendings in the simulations of the coupler, measurement errors in the waveguide geometry or width irregularities along the couplers.

As a coupler for the HOM experiment, we choose the coupler with a length of 257 μm . The splitting ratio is $(52.11 \pm 0.01) \%$ and $(57.14 \pm 0.01) \%$ depending on the input waveguides and thus close to 50:50 splitting. Here, we have to consider that the bandwidth of the coupler is reduced since we use two coupling cycles.

Hence, we found a suitable coupler for our HOM experiment with which we will continue in the next section.

HOM Interference in a TFLN Directional Coupler

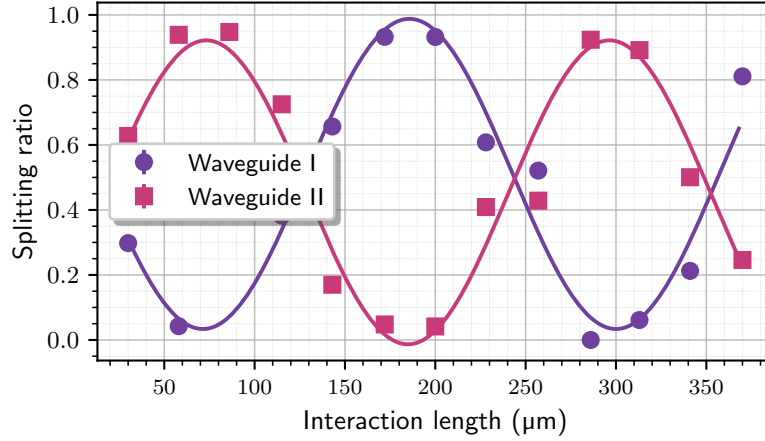


Fig. 4.5: Splitting ratio as a function of the interaction length. The light oscillates between the two waveguides. The splitting ratio is measured for both input waveguides. The error bars are so small that they are not visible in the plot.

4.4 Measuring HOM Interference in TFLN

To measure the HOM interference with a directional coupler in TFLN, we first design and build a new setup. The setup, shown in Figure 4.6, consists of two parts which are the source part and the beam splitter part.

The source part generates photon pairs via a PDC process in a periodically-poled KTP waveguide. Since the source is well established in our group and since the characterization is not the main part of this work, details about the source can be found in [126, 127]. The source is a dispersion engineered type 2 PDC source to generate photon pairs. We pump the source with a pulsed laser with a wavelength of (772.5 ± 0.5) nm and a bandwidth of around 0.3 nm which provides spectrally degenerate photon pairs. We set the mean photon number of the generated photons to below 0.01 pairs per pulse to reduce the generation of higher order photon components. Both photons from the pair have a wavelength of (1545 ± 1) nm. To guarantee spectral decorrelation between the generated photons, we place a spectral filter with a bandwidth of 1.8 nm after the sample to remove side lobes. After the filter, a polarizing beam splitter (PBS) is used to separate the two photons and both are sent through a delay line. Next, the photons are coupled to a polarization maintaining (PM) fiber to lead them to the second part of the setup.

The second part of the setup starts with the PM fibers. From here, the two photons are coupled to free space. One of the two photons is sent through a quarter and a half waveplate such that the two photons have the same polarization and are indistinguishable when impinging on the beam splitter at the same time. Then, the two photons are guided to the sample by mirrors. In one arm, we replace a mirror by a prism to bring the two photons in close proximity to be able to couple both photons with a single lens to the two input waveguides of the directional coupler. After the sample, another lens is used for outcoupling and the two photons are separated using another prism. Lastly, the two photons are guided by mirrors, each to a single-mode fiber which is connected to a superconducting nanowire single photon detector (SNSPD). The SNSPDs have a dead time of 70 ns and an efficiency of more than 95 %. The total setup efficiency including chip losses, coupling losses and detector efficiency is ~ 3 %.

4.4. Measuring HOM Interference in TFLN

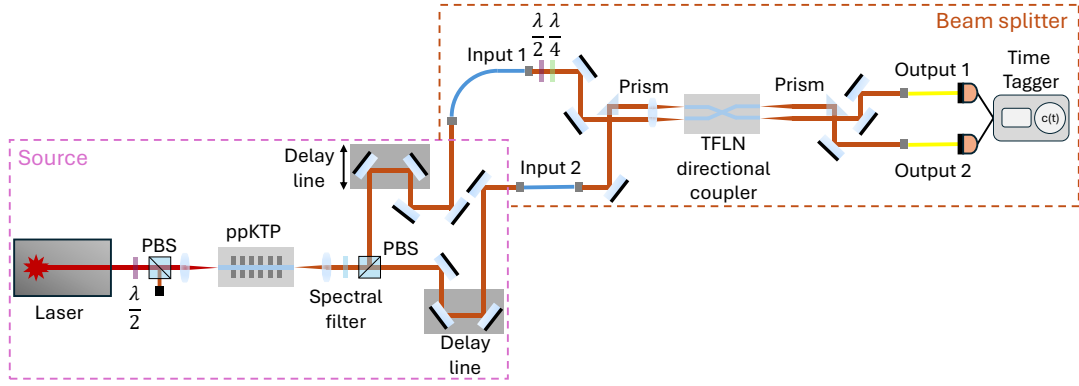


Fig. 4.6: Setup to measure the HOM dip in a TFLN directional coupler. It consists of two parts which are the source part to prepare the indistinguishable photons and the beam splitter part at which the HOM interference is measured.

To measure the HOM dip, we scan one of the delay lines. We use a step size of $1\ \mu\text{m}$ and scan a range of $250\ \mu\text{m}$. To counteract long term intensity fluctuations, we perform several scans and add them up afterwards resulting in more than 500 coincidences outside of the dip. The resulting normalized coincidences as a function of the delay are plotted in Figure 4.7. In this plot, we did not subtract background counts but normalized the baseline to 1 for better visualization. The error bars are due to click statistic. As expected, we can clearly see a dip in the coincidences when there is no time delay between the photons. At this point, the two photons are indistinguishable and leave the directional coupler in the same output waveguide since they bunch. If there is a time delay, we measure coincidences, since the two photons are distinguishable and leave the directional coupler separately. Thus, we measure the expected HOM interference behavior.

To estimate the visibility of the measured dip we apply the following Gaussian fit including

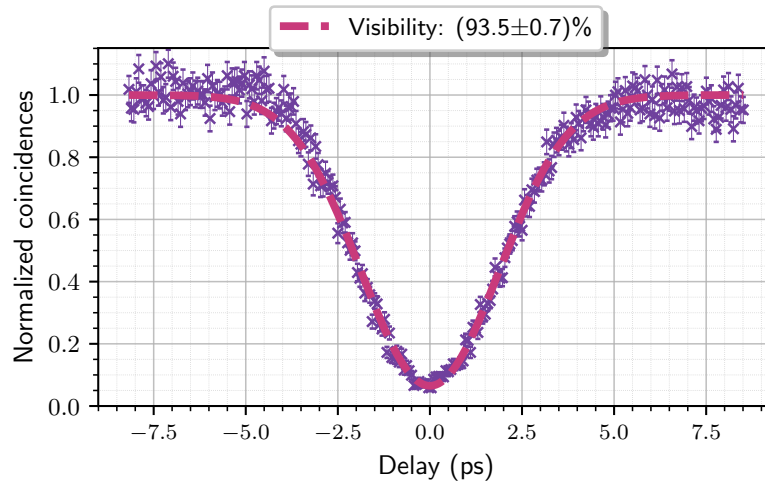


Fig. 4.7: HOM Dip in TFLN. The coincidences drop almost to zero when there is no time delay between the two photons. We measure a HOM dip with a visibility of $(93.5 \pm 0.7)\%$.

HOM Interference in a TFLN Directional Coupler

statistical error bars:

$$c(t) = 1 - de^{at^2}. \quad (4.2)$$

From this fit, we extract a visibility of $(93.5 \pm 0.7) \%$. To contextualize this value, we calculate the maximum theoretically achievable visibility which consists of the source visibility meaning the indistinguishability of the generated photons as well as the imperfect splitting ratio of the directional coupler. Thus, we first estimate the visibility of the source by using a fiber beam splitter which yields a HOM visibility of $(98.01 \pm 0.24) \%$ [128]. As mentioned above, the splitting ratio of the directional coupler was on average $(54.6 \pm 3.8) \%$. This limits the visibility to $(98.32 \pm 2.8) \%$ calculated via Equation 2.141. Hence, the theoretically maximal achievable visibility is $(96.36 \pm 2.8) \%$. Thus, we have a difference of 3 % between measurement and theory. One possible explanation for this could be that we have hybrid modes in our waveguides [129]. However, the material does not influence the HOM interference and therefore TFLN is an ideal material platform for integrated quantum optics. The demonstrated on-chip HOM interference is one of the main building blocks for integrated quantum optics and paves the way towards the integration of complex quantum circuits.

Summary

In this chapter, we showed a genuine quantum effect in TFLN. We measured the interference between two single photons in a directional coupler. For this, we designed and characterized a directional coupler. We showed a HOM dip with a visibility of $(93.5 \pm 0.7) \%$ reduced by the source visibility and the non-perfect splitting ratio of the directional coupler. The result demonstrates that we are able to design, fabricate and characterize directional couplers in TFLN and we show that the material does not effect the quantum nature of light.

5

Engineering TFLN Photon Sources

In the last chapter, we have shown a passive quantum effect in TFLN waveguides. Next, we want to expand our scope and move on to photon sources to be able to generate photons with tailored properties for specific experiments and applications. There is a variety of ways to realize photon sources such as quantum dots, color centers, metamaterials and numerous more [130]. Here, we want to focus on generating photons via a PDC process in a TFLN waveguide with tailored properties to harness the advantages of this material platform. Thus, in this chapter, we will discuss the topic of dispersion engineering of photon sources in TFLN. First, we will introduce dispersion engineering and we will show why TFLN is ideally suited for this. Then, we will investigate first order dispersion engineering in which we will match phase velocities to generate photons with a specific wavelength. Next, we will discuss how inhomogeneities in the waveguide geometry influence the phase matching function and how we can counteract those. In the end, we will move on to second order dispersion engineering by investigating the group velocities.

5.1 Dispersion Engineering in TFLN Waveguides

With dispersion engineering the characteristics of a process can be modified. For example, for a photon pair source the photons wavelength, the spectral bandwidths or the spectral correlations can be changed by engineering the dispersion of the underlying process [72]. To alter the dispersion, different methods can be used. On the one hand, we can change the material properties and on the other hand, we can change the waveguide geometry. With the latter one, we change the modal dispersion which is possible especially in high-index contrast waveguides such as TFLN. This is due to the fact that the dispersion strongly depends on the geometry of the waveguide rather than on the material dispersion. A change in the geometry leads to a modification of the effective refractive index of the TE and TM modes for various wavelengths. This change in the effective refractive index results in a variation of the group velocities and thus each waveguide geometry can have a different dispersion leading to different process characteristics. Since TFLN exhibits the possibility of dispersion engineering via the geometry, we can design three-wave mixing source with defined properties which are not achievable in titanium indiffused lithium niobate waveguides. For example, it is not possible to realize a symmetric group velocity matched PDC source at 1550 nm in titanium indiffused waveguides which is possible in other material with low-index contrast waveguides such as KTP due to its specific dispersion properties [74, 131]. However,

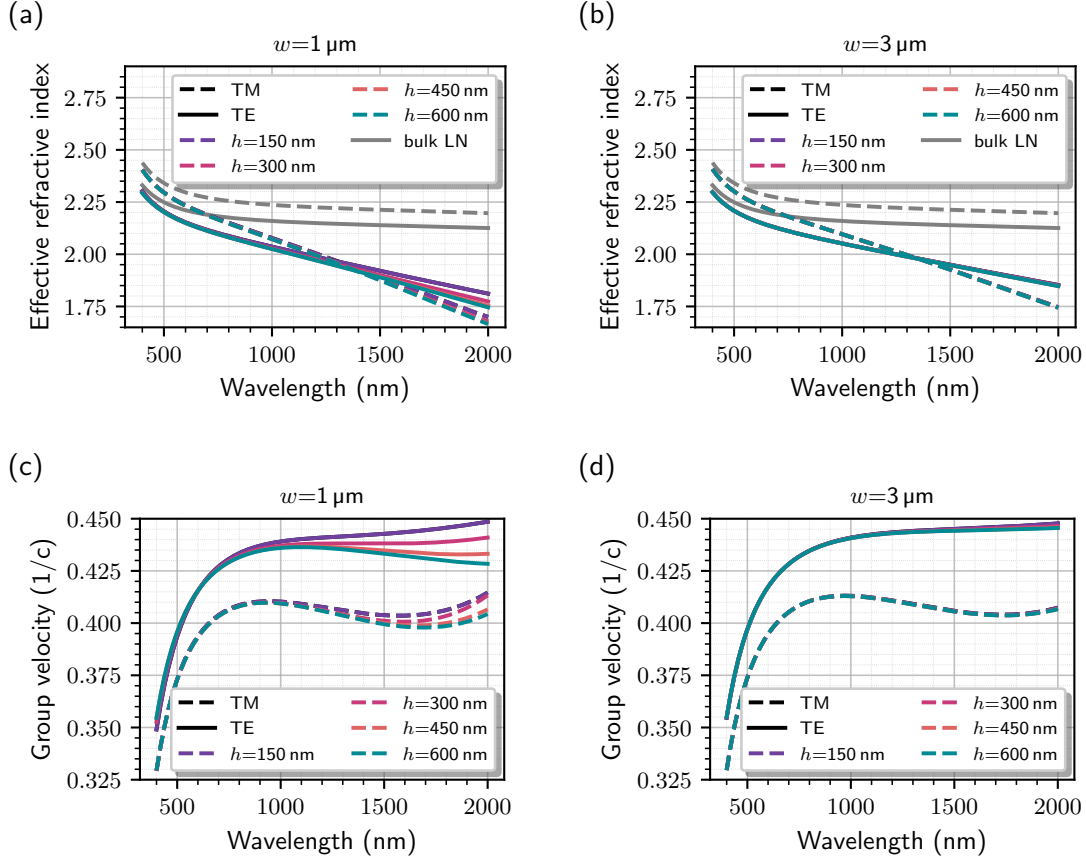


Fig. 5.1: Dispersion properties of the fundamental modes of single- and multi-mode waveguides for different etching depths. Effective refractive index as a function of the wavelength (a) for single-mode waveguides and (b) for multi-mode waveguides as well as the dispersion of bulk lithium niobate as a comparison. Group velocities as a function of the wavelength (c) for single-mode waveguides and (d) for multi-mode waveguides for which the group velocity curves overlap for different etching depths. The dispersion of a single-mode waveguide is significantly impacted by the etching depth whereas the dispersion of a multi-mode waveguide is less influenced.

with TFLN we can now achieve this specific source next to others while also harnessing the benefits of high-index contrast waveguides such as small footprint and efficient devices. This means, by tailoring the waveguide geometry we can adapt the dispersion to achieve specific processes for particular applications which has not been exploited much experimentally yet [17, 132]. However, this also indicates that fabrication tolerances can have a significant impact on the designed devices as discussed later on in this chapter.

To observe the impact of the geometry on the dispersion, we study different waveguide geometries in terms of their dispersion. We investigate the influence of the waveguide geometry via a waveguide with a narrow top width of $w=1\ \mu\text{m}$ such that the waveguide is single-mode and on the other hand, a waveguide with a broader top width of $w=3\ \mu\text{m}$ to achieve a multi-mode waveguide. The other geometry parameters are the same and as follows: thin-film thickness $t=600\ \text{nm}$, angle $a=30^\circ$, cladding $c=0\ \mu\text{m}$. To clearly see the influence of the geometry on the dispersion, we vary the etching depth from $h=150\ \text{nm}$

5.1. Dispersion Engineering in TFLN Waveguides

(shallow etched) to $h=600$ nm (completely etched) in 150 nm steps. For these waveguide geometries, we simulate the effective refractive index and calculate the group velocities as a function of the wavelength. Moreover, we also plot the refractive index of bulk lithium niobate in grey as a comparison to the geometry modified dispersion of TFLN waveguides. The results are plotted in Figure 5.1. In panel (a) and (b), we plot the effective refractive index as a function of the wavelength for single-mode and multi-mode waveguides, respectively. The effective refractive index of the TE mode is shown as a solid line and of the TM mode as a dashed line. For each etching depth we use a different color. When comparing the dispersion of bulk lithium niobate to the dispersion of the TFLN waveguides, we can see that the dispersion of the TFLN waveguides differs significantly from the dispersion of bulk lithium niobate. Both have at a wavelength of 400 nm almost the same (effective) refractive index. After that, the dispersion of the TFLN waveguides decreases faster than that of bulk lithium niobate. In addition, for the TFLN waveguides the effective refractive index of the TM mode reduces more than that of the TE mode which leads to a crossing of the two curves at around 1300 nm. This crossing means that for one wavelength the TE and TM mode have the same effective refractive index. For bulk lithium niobate such a crossing does not occur and the refractive index of the TM mode is always higher than that of the TE mode. This significant difference in the dispersion between bulk lithium niobate and TFLN waveguides demonstrates that the dispersion of TFLN waveguides is more impacted by the waveguide geometry than by the material dispersion due to the higher confinement. Next, we want to exemplarily investigate the impact of the etching depth on the dispersion to show how much influence the geometry has on the dispersion. We can observe that the etching depth changes significantly the effective refractive index of a single-mode waveguide for the TE and TM mode especially in the telecom regime, see panel (a). For a multi-mode waveguide (see panel (b)), the effective refractive index does not change significantly with the etching depth. This different behavior is due to the confinement of the mode in the waveguide. For a broader waveguide, the mode is less confined and thus less influenced by the waveguide geometry. For a narrow waveguide, the mode has a higher confinement and thus, a change in the etching depth leads to a change of the effective refractive index of the mode. This indicates that a single-mode waveguide is more influenced by changes in the geometry during fabrication than a multi-mode waveguide. Therefore, when designing a waveguide for a specific purpose, one has to take fabrication tolerances into account since these have an influence on the geometry and change the dispersion.

To investigate this further, we plot the group velocities as a function of the wavelength, see Figure 5.1 (c) and (d). If we compare the behavior of the group velocity of the differently deep etched single-mode waveguides (see Figure 5.1 (c)), we can see a change of the group velocity with the etching depth. Especially, in the telecom region, the group velocity changes significantly. By etching deeper, we get a lower group velocity. Furthermore, also the slope changes which is important for dispersion engineered sources as discussed in section 2.5.1. For the TE mode, we go from a positive slope to a negative slope if we etch completely through. For the TM mode, the slope is not altered, but the group velocity also decreases with increasing etching depth. The comparatively less significant change in the TM mode is due to that fact that the TM mode is polarized along the vertical direction. This means, the mode is more altered by changes in the vertical direction as for example the thin-film thickness compared to horizontal changes as the etching depth of the waveguide.

Moreover, we can observe that the change in the group velocities in the telecom region for the TE mode from one etching depth to the next is the same except for $h=450$ nm to

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$h=600$ nm. The change between the group velocity of an etching depth of 450 nm and the group velocity of an etching depth of 600 nm is smaller compared 300 nm to 450 nm etching depth. For the TM mode, we can see that the change in the group velocity between 150 nm etching depth and 300 nm etching depth is larger compared to the other etching depths. This indicates that some geometries are less influenced by a specific geometry parameter and thus more tolerant to fabrication imperfections than others [133]. To achieve a stable geometry, it is crucial that the waveguide mode is not altered significantly by slight changes in the geometry. This is for example the case for a deep etched waveguide. For this configuration, the mode is confined in the etched waveguide and does not reach significantly into the remaining slab. Thus, if the etching depth is changed slightly, the mode will still be confined in the etched waveguide and will not change substantially. This is analog to a broad waveguide top width. However, if the mode reaches into the slab for example for a not deep etched waveguide, an increase of the etching depth will lead to the fact that the mode will be shifted from the slab to the etched waveguide which will change the mode and thus also the effective refractive index. For the top width, this is similar to a small waveguide width.

If we compare the narrower waveguide to the broader waveguide (Figure 5.1 (d)), we can see that for the broader waveguides a change in the etching depth does not have the same impact as for a narrower waveguides (the group velocities from all etching depths are almost completely on top of each other).

The results clearly indicate that in TFLN waveguide geometry based dispersion engineering is possible since we can not only change the value of the effective refractive index but we can also alter the slope of the group velocities which are important for designing different kinds of PDC sources as discussed in section 2.5.1. To conclude, this shows that TFLN is ideally suited for dispersion engineering since the dispersion can strongly depend on the waveguide geometry due to the high mode confinement.

In this thesis, we will split dispersion engineering into two types. The first investigation depth is to simply match the phase velocities to generate photons at desired wavelengths. In the second one we extend our investigation to group velocities to also achieve a specific phase matching function.

5.2 First-Order Dispersion Engineering

In first-order dispersion engineering, we are matching phase velocities to achieve an efficient PDC process at desired wavelengths. This means, we omit the phase matching angle or correlations in the JSA. This is achieved by phase velocity matching via periodic poling (see section 2.3.2). We target a PDC source generating photons at distinct wavelengths. To motivate this, we first discuss the state of the art of single pass periodically poled PDC sources realized at room temperature in TFLN waveguides in terms of wavelengths.

5.2.1 State of the Art - First-Order Dispersion Engineering

The first PDC source in TFLN was published in 2019 by Chen et al. [19]. They realized a type 0 PDC source pumped at 775 nm to generate photons at 1550 nm. After this, various more PDC sources have been published. Most of the sources are type 0 sources pumped at 775 nm to generate photons at 1550 nm [28, 29, 30, 134, 135, 136, 124]. In 2022, a type 2

5.2. First-Order Dispersion Engineering

source was demonstrated by Xin et al. with the same wavelength combination of 775 nm and 1550 nm [137]. However, also PDC sources with other wavelength combinations were demonstrated. In 2024, Harper et al. published a type 0 PDC source pumped at 400 nm to create light at 800 nm [138]. In the same year, Williams et al. showed a type 0 PDC source from 1 μm to 2 μm [139].

When examining the wavelength combinations between the pump photon and the signal and idler photons, we can see that all above mentioned sources were spectrally degenerate meaning the signal and idler photon have the same wavelength. However, a highly non-degenerate PDC source with one photon in the visible wavelength regime and the other photon in the infrared is a crucial building block for various applications in integrated quantum optics. It can be utilized for example as a cost-effective, highly efficient heralded single photon source for quantum communication schemes [140, 141]. A low-cost silicon avalanche single photon detector can be used to measure the photon with the shorter wavelength to herald the existence of the other photon. Moreover, such a source can be used in the field of quantum spectroscopy for measurement schemes with undetected photons [142, 143, 144, 145, 146]. Here, the longer wavelength is used to interact with an object under test while the shorter wavelength is measured. With this, we can use again the shorter wavelength for measuring while gaining knowledge about an object's properties in the infrared.

Thus, in the next part we will describe how we engineer a highly non-degenerate PDC source in first order. We start with simulations of the effective refractive index of waveguide geometries and discuss the influence of fabrication tolerances. Afterwards, we explain the optical setup used for source characterization and experimental procedures. Then, we first investigate the spectral characteristics of the source and afterwards perform photon counting experiments to evaluate the quantum properties. We also show that the waveguide simulations match very well with the experiment.

The following results were published in [147].

5.2.2 Design of a Non-Degenerate Photon-Pair Source

As mentioned above, we target to demonstrate a highly non-degenerate PDC source in a TFLN waveguide. We aim to generate the photons at 810 nm (signal) and 1550 nm (idler) meaning pumping at 532 nm in a type 0 process to utilize the highest tensor element d_{33} to achieve an efficient source. First, we set our fabrication limitations and tolerances. At

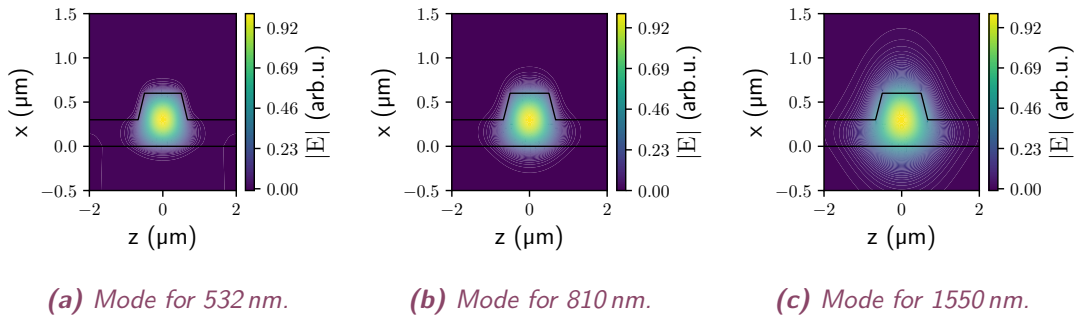


Fig. 5.2: Simulated mode distribution of a cross-section of a waveguide for the non-degenerate PDC design.

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that time, the thin-film thickness was fixed to 600 nm of 5 mol% MgO doped x-cut lithium niobate and the angle was fixed due to the etching process to 30° . We used an etching depth of 300 nm since this etching depth could be reliably fabricated. Moreover, to protect the sample, we applied a $1\ \mu\text{m}$ thick SiO_2 cladding layer.

For an efficient process, we design the waveguide to be single-mode at 1550 nm to suppress higher order processes. To find a single-mode waveguide geometry, we simulate the modes of different waveguide top widths with the *MODE* solver of *Lumerical* and the model explained in section 2.4.2. From this analysis, we choose a waveguide with a top width of $1\ \mu\text{m}$. In Figure 5.2, we plot the fundamental TE modes of this waveguide geometry for the three used wavelengths.

After setting the waveguide geometry, in the first-order dispersion engineering the only remaining constraint is the momentum conservation which we achieve by periodic poling (see section 2.3.2). However, we know from the investigation of the group velocities on the geometry above that the dispersion of TFLN waveguides strongly depends on the waveguide geometry due to the high mode confinement. This results in the fact that a slightly differently fabricated waveguide geometry would require a different poling period to achieve phase matching at the same wavelengths as the initial geometry. Therefore, it is crucial to know each geometry parameter as precisely as possible and to take fabrication tolerances into account when designing a sample to be able to generate photons at the desired wavelengths. One parameter that can be measured before the waveguide fabrication that has a vital impact on the dispersion is the thin-film thickness [133]. Thus, we measure the thin-film thickness of the $1\text{ cm} \times 1\text{ cm}$ sample at five different points in a $7.5\text{ mm} \times 7.5\text{ mm}$ area with an ellipsometer. We average these five points and calculate the standard deviation which results in a thin-film thickness of $(607 \pm 1)\text{ nm}$. With this thin-film thickness, we calculate with Equation 2.116 a poling period of $3.314\ \mu\text{m}$ for the ideal waveguide geometry ($t=607\text{ nm}$, $h=300\text{ nm}$, $w=1\ \mu\text{m}$, $a=30^\circ$, $c=1\ \mu\text{m}$).

However, as mentioned above the fabricated waveguides can due to the fabrication process slightly differ from the simulated one. This leads to a different dispersion which alters the necessary poling period to achieve phase matching at the desired wavelengths. Thus, it is important to factor in fabrication tolerances when designing a sample to reach the intended process. To take possible fabrication tolerances into account, we use the method proposed in [17]. In this method, we investigate the influence on the propagation constants of each geometry parameter individually to estimate a poling period range based on the expected fabrication tolerances. To estimate for example the influence of the etching depth, we simulate the effective refractive index for etching depths in the range of 200 nm to 400 nm in 50 nm steps. For each geometry, we calculate with Equation 2.116 the phase mismatch and take as a poling period $\Lambda = 3.314\ \mu\text{m}$ from the ideal geometry. We plot the phase mismatch as a function of the etching depth and use a linear fit to determine the slope $\frac{\delta\Delta\beta}{\delta h}$. This procedure is done for all geometry parameters and the plots are shown in Figure 5.3. When investigating the slopes of Figure 5.3 it can be seen that only the slope for the etching depth is positive and that all other slopes are negative. In general, the slopes of the individual

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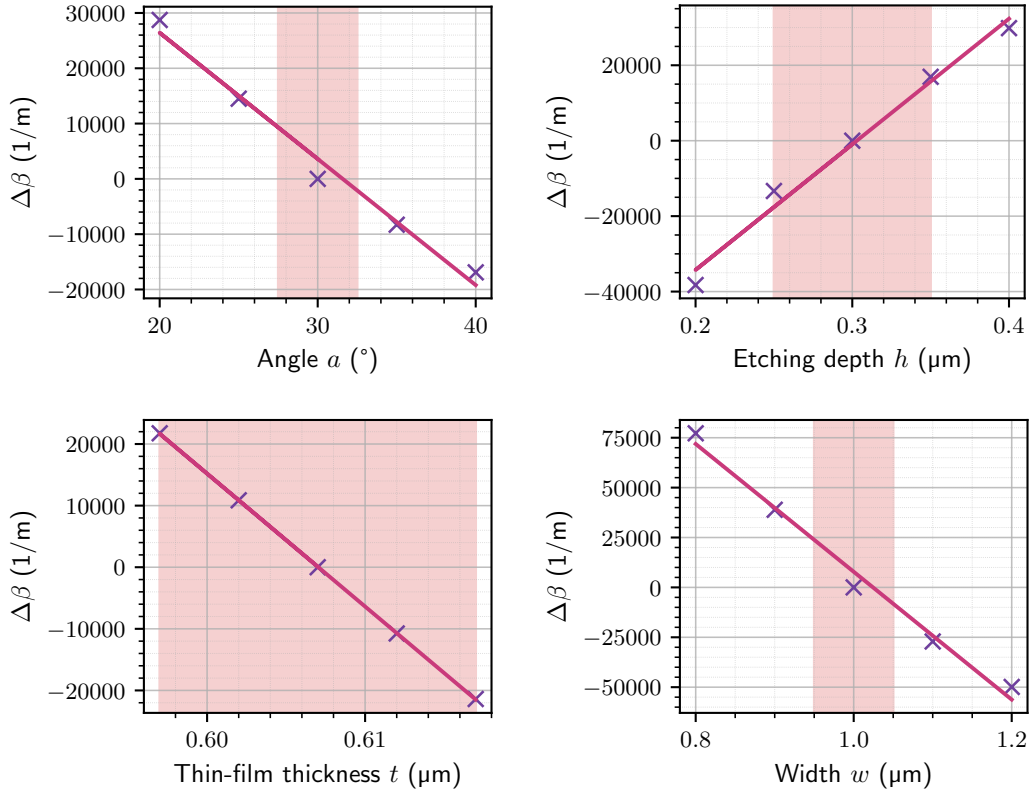


Fig. 5.3: Phase mismatch as a function of each individual waveguide geometry parameter. We apply a linear fit to obtain the slopes. Moreover, we mark in pink the fabrication tolerances.

parameters are as follows:

$$\begin{aligned}\frac{\delta\Delta\beta}{\delta a} &= -2.3 \cdot 10^3 \frac{1}{m^\circ} \\ \frac{\delta\Delta\beta}{\delta h} &= 332.9 \cdot 10^3 \frac{1}{m \mu m} \\ \frac{\delta\Delta\beta}{\delta t} &= -2.2 \cdot 10^6 \frac{1}{m \mu m} \\ \frac{\delta\Delta\beta}{\delta w} &= -320.2 \cdot 10^3 \frac{1}{m \mu m}.\end{aligned}$$

Here, it can be seen that the slope of the thin-film thickness has the highest value and thus the highest impact. However, this is a parameter mainly given by the manufacturer. Thus, later in this thesis, we will discuss the influence of the thin-film thickness variations as well as approaches to counteract those. With this, we can calculate the phase mismatch due to the fabrication tolerances [17]:

$$\Delta\beta_{\text{sum}} = \frac{\delta(\Delta\beta)}{\delta w} \cdot \Delta w + \frac{\delta(\Delta\beta)}{\delta h} \cdot \Delta h + \frac{\delta(\Delta\beta)}{\delta a} \cdot \Delta a + \frac{\delta(\Delta\beta)}{\delta t} \cdot \Delta t. \quad (5.1)$$

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Here, Δw , Δh , Δa and Δt are the fabrication tolerances of each parameter. We evaluate the fabrication tolerances by measuring the waveguide geometry on multiple samples and calculate the standard deviation. We use an atomic force microscope to measure the sidewall angle as well as the top width. To estimate the fabrication tolerance for the etching depth, we use an atomic force microscope as well as ellipsometer. This leads us to the following fabrication tolerances $\Delta w = 100 \text{ nm}$, $\Delta h = 100 \text{ nm}$, $\Delta a = 5^\circ$ and $\Delta t = 20 \text{ nm}$ which we marked as pink regions in Figure 5.3. These tolerances can lead to a phase mismatch of $-54230 \frac{1}{m}$ resulting in a wavelength shift for the idler photon of a 100 nm at worst. Thus, we want to estimate a poling period range to be able to counteract this. This is done via the following equation in which we add/subtract the phase mismatch due to the fabrication tolerances from the phase mismatch of the initial geometry $\Delta\beta_0$ to calculate the minimal and maximal poling period:

$$\Delta\Lambda = \frac{2 \cdot \pi}{\Delta\beta_0 \pm \Delta\beta_{\text{sum}}}. \quad (5.2)$$

This leads to a poling period range between $3.266 \mu\text{m}$ - $3.443 \mu\text{m}$ with which we should be able to generate photons with the desired wavelengths even if fabrication errors occur. Thus, on our sample we put the following poling periods: $3.266 \mu\text{m}$, $3.310 \mu\text{m}$, $3.353 \mu\text{m}$, $3.397 \mu\text{m}$ and $3.443 \mu\text{m}$. Based on this, we fabricate a periodically poled waveguide sample. To characterize the sample, we use the following setup.

5.2.3 Experimental Setup

To characterize the non-degenerate PDC source, we use the setup shown in Figure 5.4. It consists of the main source part and three different measurement parts.

The main part displayed in (a) starts with three different lasers. The first laser is used as a pump laser for the PDC source and it is a tunable CW-laser used at 534 nm. The second laser is a pulsed laser with central wavelength of 810 nm generating ultrashort pulses at repetition rate of 80 MHz. The third laser is a tunable CW-laser from 1490 nm to 1640 nm. The first laser is sent through a half-wave plate and a polarizing beam splitter to set the polarization and to control the power. The other two laser are also sent through half-wave plates for polarization control. Then, the three lasers are combined to one main beam line using dichroic mirrors. A single aspheric lens is used to couple the light to the periodically poled TFLN sample which is placed on a home-build oven to stabilize the temperature. After the sample, the light is collected again using a single aspheric lens. The three wavelengths are split by dichroic mirrors and coupled into separate optical fibers due to which we can easily switch between measurement devices to perform various experiments.

The first measurement setup (see Figure 5.4 (b)) is used to measure the PDC spectra of the signal and idler photons. For this, the sample is pumped with the green laser with the two other laser switched off and the signal and idler photons are measured with single-photon sensitive spectrometers. The spectrometer for the signal photons consists of a spectrograph combined with a CCD camera and for the idler photon of a spectrograph with a iDus InGaAs camera.

To map the phase matching function, a second measurement setup is utilized which is shown in Figure 5.4 (c). Here, we use the pulsed laser at 810 nm and the CW-laser in the telecom range and measure the SFG spectrum with a single-photon sensitive spectrometer consisting of a spectrograph combined with a CCD camera.

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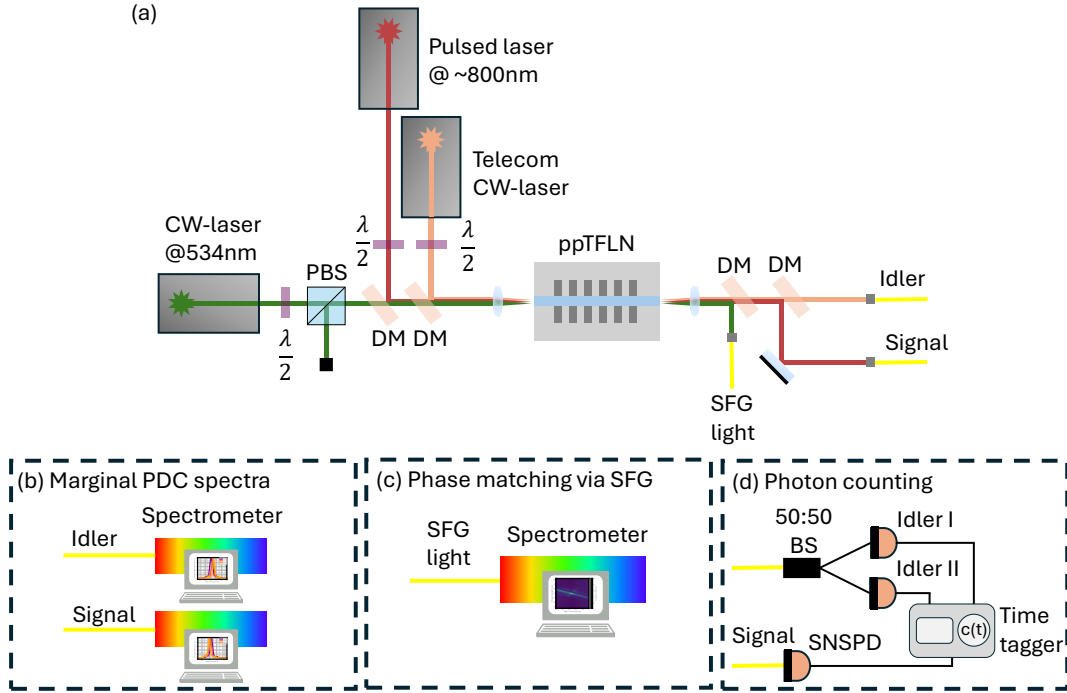


Fig. 5.4: Setup to measure a highly non-degenerate PDC source in TFLN. It consists of the main part which includes the periodically poled TFLN sample shown in (a) and three measurement setups. (b) Setup to measure the PDC spectra of the signal and idler photons with single-photon sensitive spectrometers. (c) Measurement of the phase matching function via SFG. (d) Photon counting measurement setup to evaluate the photon number purity properties.

With the last measurement setup (see Figure 5.4 (d)), we perform photon counting experiments. We introduce a fiber coupled 50:50 beam splitter in the idler arm and measure the two outputs as well as the signal photons with SNSPDs. With this setup, we characterize the PDC source which will be explained in the next section.

5.2.4 Source Characterization

In this section, we will analyze the spectral properties of the generated photons and compare the results to our simulations. Afterwards, we will perform photon counting experiments to benchmark the quantum properties of the source.

Spectral Characterization

To characterize the spectral properties of our source, we measure the dependency of the generated output wavelengths on the poling period. For this, we use the setup shown in Figure 5.4 (a) and (b). We pump our source with a slightly different pump wavelength of 534 nm instead of 532 nm due to a minor difference in the fabricated waveguide geometry which alters the phase matching condition. To find the proper pump wavelength, we measure the output spectrum for different pump wavelengths and study the signal strength to find the pump wavelength which gives the highest signal.

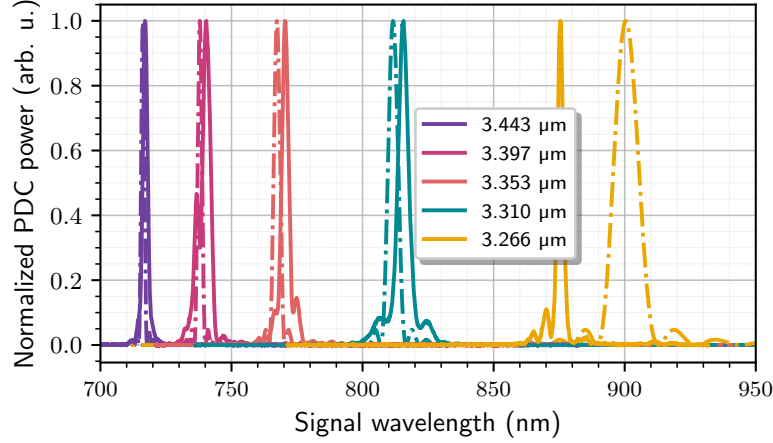


Fig. 5.5: PDC spectra of the signal photons normalized to each maximum peak power of five waveguides with different poling periods. The signal wavelength decreases with increasing poling period as expected. Moreover, we observe a good agreement between the measurements (solid lines) and the simulations (dashed lines). For the yellow spectrum the simulations does not match due to a cut-off wavelength of a dichroic mirror in the setup.

We measure the generated signal and idler photons with single-photon sensitive spectrometers. In Figure 5.5, we plot the normalized PDC power of the signal photons (solid lines) as a function of the wavelength of five waveguides with the same geometry but different poling periods. Furthermore, we also plot our simulated spectra (dashed lines). For that, we calculate the phase matching function for each poling period thereby considering varying phase matching angles. We simulate the JSI with a pump function centered at the wavelength used in the experiment and plot the projection onto the signal axis. The simulations as well as the experiment show that the signal wavelength decreases with increasing poling period. The measured trend as well as the measured peak positions are in good agreement with the simulations. A minor offset between the measured and simulated peak positions can be due to a slightly different waveguide geometry, alignment or chromatic aberrations. However, there is a clear difference between the experiment and the simulation for the waveguide with a poling period of $3.266 \mu\text{m}$ (yellow spectra). This is due to the fact that the dichroic mirror which is used to reflect the signal photons towards the fiber coupling has a cut-off at a wavelength of 900 nm . This means, we did not measure the complete spectrum but only a part of it which was reflected by the dichroic mirror. This also explains the missing side lobes at the right of the spectrum since the spectrum was cut at around 880 nm . Using a different dichroic mirror with a higher cut-off wavelength could lead to measured spectrum closer to the simulation.

Comparing the other four spectra, we can observe that the shape of each spectra is to some degree different. For example, the waveguide with a poling period of $3.310 \mu\text{m}$ (teal spectra at around 815 nm) shows stronger side peaks compared to the other three spectra. Moreover, the bandwidth of the waveguide with a poling period of $3.397 \mu\text{m}$ (magenta spectra) is broader compared to the neighboring spectra. This behavior can be due to a varying thin-film thickness which will be discussed in a later stage of this chapter.

With this source, we aimed to generate photons with a specific wavelength combination of 810 nm and 1550 nm . This is achieved with the waveguide with a poling period of $3.31 \mu\text{m}$

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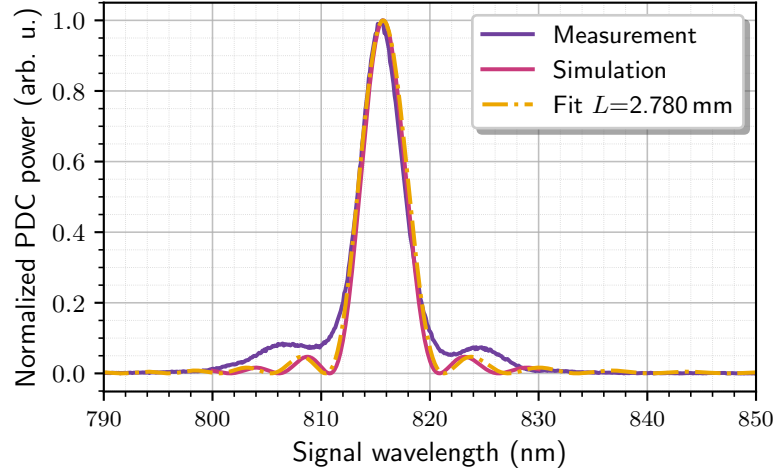


Fig. 5.6: Normalized PDC spectrum of the signal photons as a function of the signal wavelength to determine the effective poling length. In purple, we show the measured spectrum. The simulation is plotted in magenta and the applied fit function in yellow.

(teal spectra). This waveguide allows for generating photons at 815 nm and 1550 nm. Therefore, we investigate the shape and bandwidth of this spectrum in more detail. For this, we plot the measured phase matching spectrum of this specific waveguide in Figure 5.6 in purple. A precise control of the shape of the phase matching function is crucial, since a distorted phase matching function can lead to a change of the correlations in the JSI. One aspect of the shape of the spectrum is the spectral bandwidth. The spectral bandwidth is dependent on the length of the poling structure for a fixed dispersion and a longer poling structure leads to a narrower spectral bandwidth. A broadening of the bandwidth can occur either due to varying geometry parameters along the waveguide or inhomogeneities in the poling [148, 149]. This is due to the fact that the phase matching condition is fulfilled for slightly different wavelengths if the geometry or the poling period changes. Thus, it is crucial to compare the simulated and measured bandwidth to quantify the poling quality as well as the waveguide homogeneity. For this, we calculate the effective poling length. We fit to our measurement a sinc^2 -function, as discussed in section 2.3.2, which relies on the simulated dispersion data (yellow line in Figure 5.6). Moreover, we also plot again the simulated phase matching spectra in magenta. Here, we optimize the overlap of the simulation and measurement by varying the poling period to $3.305 \mu\text{m}$ achieving the same central wavelength. We can see that the measured spectrum mirrors a near to perfect sinc^2 . Evaluating the sinc^2 -fit leads to an effective length of 2.78 mm which is 93 % of our nominal length of 3 mm. This shows our excellent fabrication quality of waveguides as well as our high poling quality. The minor broadening can be attributed to long range noise in the thin-film thickness along the waveguide [149]. This noise leads to a change in the dispersion which results in the fact that the phase matching condition is fulfilled for altered wavelengths which broadens the spectrum [150]. The main impact on the shape of the phase matching function has a varying thin-film thickness [133]. Thus, we will discuss this effect as well as methods to counteract those variations later in section 5.3.

Up to this point, we can conclude that we are able to design a PDC source such that it generates photons at our desired wavelengths. We have an excellent agreement between simulation and experiment. Moreover, the measurements indicate a high quality fabrication

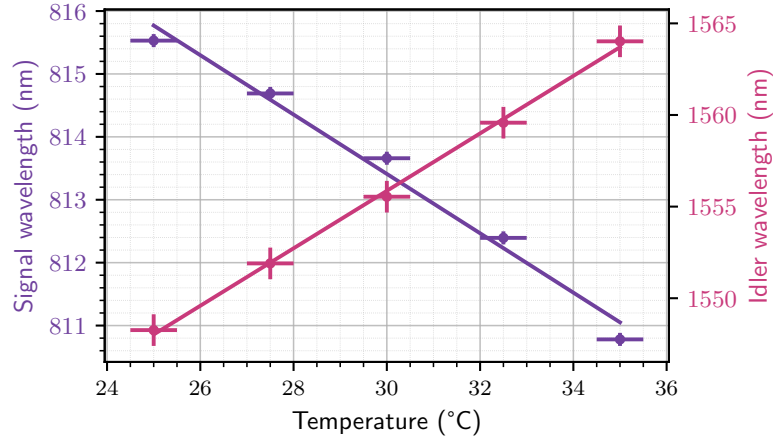


Fig. 5.7: Temperature dependency of the signal and idler photons. For an increasing temperature the signal wavelength decreases while simultaneously the idler wavelength increases.

of the waveguides and the poling.

As a next step, we want to investigate the temperature dependency of the generated output wavelength to characterize the tunability of our source. For this, we measure the spectrum of the signal and idler photons generated in a waveguide with a poling period of $3.310\ \mu\text{m}$ by using a single-photon spectrometer. We realize the tunability by varying its temperature from $25\ ^\circ\text{C}$ to $35\ ^\circ\text{C}$ in $2.5\ ^\circ\text{C}$ steps. Then, we estimate the peak wavelength position for each temperature which are shown in Figure 5.7 for the signal and idler photons as a function of the temperature. The errors bars for the wavelength are determined by the resolution of the spectrometer ($0.1\ \text{nm}$ for the signal wavelength and $0.86\ \text{nm}$ for the idler wavelength) and the accuracy of the temperature controller ($0.5\ ^\circ\text{C}$). The plot shows a decrease of the signal wavelength with increasing temperature and increase of the idler wavelength with increasing temperature. Via a linear fit, we evaluate a temperature dependency of the signal wavelength of $-0.472\ \text{nm/K}$ and for the idler wavelength of $1.57\ \text{nm/K}$. From simulations we get marginally different temperature dependencies of $-0.61\ \text{nm/K}$ for the signal wavelength and $2.28\ \text{nm/K}$ for the idler wavelength. The linear trend is the same and we attribute the deviations to minor inaccuracies in the simulated dispersion data as well as a difference in the set temperature of the heater and the actual temperature of the sample. Nonetheless, this temperature dependency shows a great tunability of our source which makes it versatile for various applications with specific wavelength requirements.

To further characterize the source in depth, we want to map the 2D phase matching function of our source to compare the measured spectrum with our simulations in terms of the phase matching angle. To characterize the phase matching function, we use the setup shown in Figure 5.4 (a) and (c). We utilize the SFG process to map the phase matching function and pump the source with the pulsed laser at around $815\ \text{nm}$ and the CW-laser in the telecom range. To measure the phase matching function, we set the pulsed laser at a specific wavelength and scan the CW-laser from $1500\ \text{nm}$ to $1600\ \text{nm}$ in $0.2\ \text{nm}$ steps. For each wavelength of the CW-laser, we measure with a single-photon sensitive spectrometer the spectrum of the generated SFG light. Afterwards, we change the wavelength of the pulsed laser by $5\ \text{nm}$ and repeat the wavelength scan of the telecom laser. We scan the pulsed laser from $795\ \text{nm}$ to $825\ \text{nm}$ in $5\ \text{nm}$ steps. Next, we combine all measured spectra. Lastly, we

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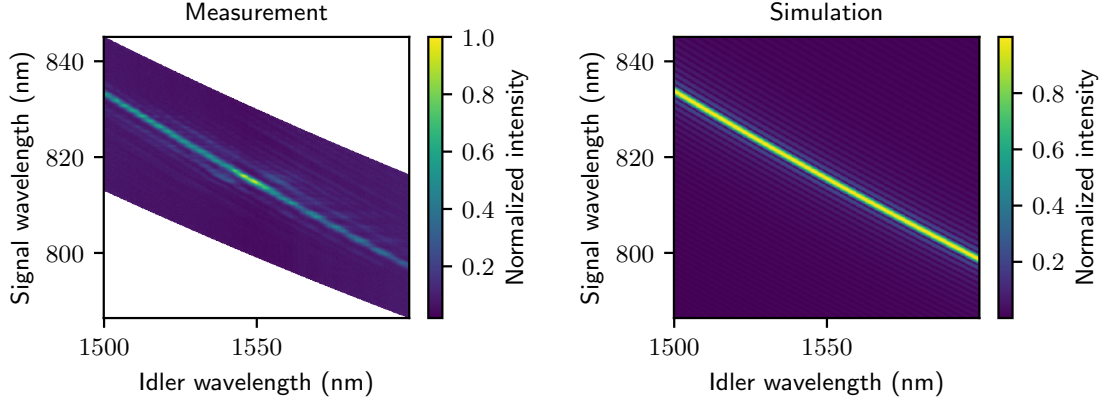


Fig. 5.8: Measured and simulated phase matching function. The measured phase matching function exhibits a phase matching angle of $(-52.60 \pm 0.11)^\circ$ which agrees excellently with the simulated angle of -51.66° .

translate the measured wavelength from the spectrometer via energy conservation to the signal wavelength to retrieve a signal and idler plot for better visual comparison between measurement and simulation. Both plots are shown Figure 5.8. In the measurement, we can observe a brighter spot at 1550 nm and 815 nm which is due to the fact that the setup was aligned for this wavelengths and that the coupling is wavelength dependent. For this point, we measure around 3500 counts in 1s with dark counts of around 50 counts. Comparing the measurement to our simulation, we can observe an excellent agreement. The measured phase matching angle of $(-52.60 \pm 0.11)^\circ$ slightly deviates from the simulated one of -51.66° which is mostly due to a minor difference between the fabricated and simulated waveguide geometry. This verifies that our simulations of the dispersion of TFLN waveguides are perfectly suited for the design of tailored photon pair sources with specific wavelength requirements.

Photon Number Purity

To use the source in any quantum optics experiment such as quantum communication, it is crucial to know the number of photon pairs generated in the source. Thus, to benchmark the quantum performance of our source, we perform photon counting experiments and evaluate the heralded second-order correlation function to determine the photon number purity. For this, we use the setup as shown in Figure 5.4 (a) and (d). We measure the single count rates of the signal photons R_s , the idler photons R_1 and R_2 , the twofold coincidences between signal and idler 1 $R_{s,1}$ and signal and idler 2 $R_{s,2}$ as well as the threefold coincidences $R_{s,1,2}$ as a function of the pump power. To monitor the pump power, we use a power meter placed after the sample in front of the fiber coupling for the SFG light. In Figure 5.9 (a) the count rates as a function of the pump power are shown. Here, we do not subtract any accidental counts from the two- and threefold coincidences due to the high coincidence-to-accidental ratio (CAR) which is shown in Figure 5.9 (b). For the highest pump power of $5.7 \mu\text{W}$ the CAR is 12 and increases to 1008 for the lowest pump power of 65.8 nW . The error bars are instrument-limited for the pump power and Poissonian noise for the count rates. The fact that we are measuring twofold coincidences shows that the PDC source generates correlated photon pairs. First, we evaluate the Klyshko efficiencies from the count rates to

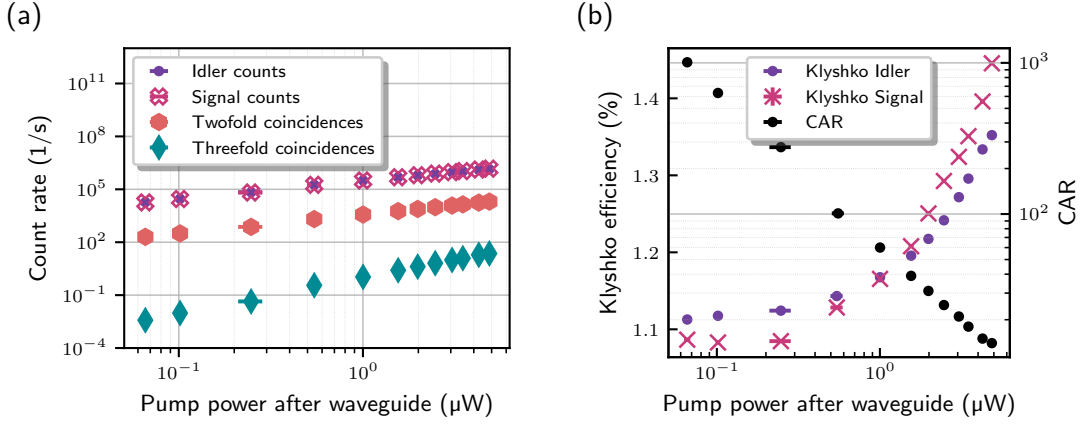


Fig. 5.9: (a) Count rates as a function of the pump power. (b) Klyshko efficiency of the signal and idler photons as well as the coincidence-to-accidental ratio as a function of the pump power. The error bars are so small that they are not always visible.

be able to later re-calculate the actual generated photon pairs inside the source. Since we introduce a 50:50 beam splitter in the idler path, we can not use the equation as discussed in section 2.5.1, but have to adjust them slightly:

$$\eta_s = \frac{R_{s,1} + R_{s,2}}{R_1 + R_2} \quad (5.3)$$

$$\eta_i = \frac{R_{s,1} + R_{s,2}}{R_s}. \quad (5.4)$$

In Figure 5.9 (b) the Klyshko efficiencies for the signal and idler photons are displayed as a function of the pump power. For low pump powers the Klyshko efficiencies stay constant and both increase for higher pump powers. This increase is due to the generation of higher-order photon number components. However, the Klyshko efficiencies increase differently for the signal and idler arm. This indicates different losses for the two wavelengths inside the waveguide as well as different coupling efficiencies of the signal and idler photons.

For low pump powers the Klyshko efficiency of the signal photon is constant at 1.08 % and at 1.1 % for the idler photon. These efficiencies are rather low, but can be significantly improved by improving the in- and out-coupling efficiencies by using waveguide tapers or grating couplers [151, 152, 153, 154].

As mentioned above, the increase of the Klyshko efficiency at higher pump power is due to the generation of multiple photon pairs. This is also indicated by the increase of the two- and threefold coincidences (see Figure 5.9 (a)). To quantify the generation of single and multiple photon pairs, we calculate the heralded second-order correlation function at zero-delay $g_h^{(2)}(0)$. Furthermore, we also want to compare our measurements with the theory. To calculate the heralded second-order correlation function theoretically we use the following equation [155]

$$g_h^{(2)}(0, n) = \frac{\sum_{\tilde{n}} \tilde{n}^2 (\tilde{n} - 1) \rho(\tilde{n})_{\text{PDC}}}{[\sum_{\tilde{n}} \tilde{n}^2 \rho(\tilde{n})_{\text{PDC}}]^2} \sum_{\tilde{n}} \tilde{n} \rho(\tilde{n})_{\text{PDC}} \quad (5.5)$$

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where n is the mean photon number per coincidence window and $\tilde{n}$ the photon number running from zero to infinity. Moreover, we assume for $\rho(\tilde{n})_{\text{PDC}}$ Poissonian statistics [156]:

$$\rho(\tilde{n})_{\text{PDC}} = \frac{n^{\tilde{n}}}{\tilde{n}!} \cdot e^{-n} \quad (5.6)$$

due to the great number of spectral modes in the source. Since the theoretical expression for the $g_h^{(2)}(0)$ is determined by the mean photon number, we first have to calculate the measured mean photon number to be able to compare both. To do so, we multiply the PDC rate by the coincidence window of the time tagger which we set to 1 ns. In Figure 5.10 (a), we plot the mean photon number as a function of the pump power. The plot shows that the mean photon number increases from $1.66 \cdot 10^{-3}$ to $103 \cdot 10^{-3}$. We apply a linear fit which shows that for high pump powers the mean photon number starts to deviate from the linear trend. We assign this behavior to detector saturation effects. Next, we calculate the heralded second-order correlation function $g_h^{(2)}(0)$ via Equation 2.132 and plot the results as a function of the mean photon number together with the theory (see Figure 5.10 (b)). We can observe a $g_h^{(2)}(0)$ value for low pump powers as low as $(6.7 \pm 1.1) \cdot 10^{-3}$ and an increase of the $g_h^{(2)}(0)$ value for higher pump powers due to the higher probability of generating multiple photon pairs. This resembles the expected behavior and shows that our source is capable of generating high quality heralded single photons. Moreover, when comparing the measurements with the theory we can see an excellent agreement. This demonstrates that we are able to extract high quality heralded single photons from our source. As a last characterization step, we will evaluate the brightness of our source.

Brightness

To determine the brightness of our source, we utilize Equation 2.133. The brightness depends on the pump power inside the waveguide. Since we measured the pump power after the waveguide, we first have to estimate from this the pump power inside of it. To do

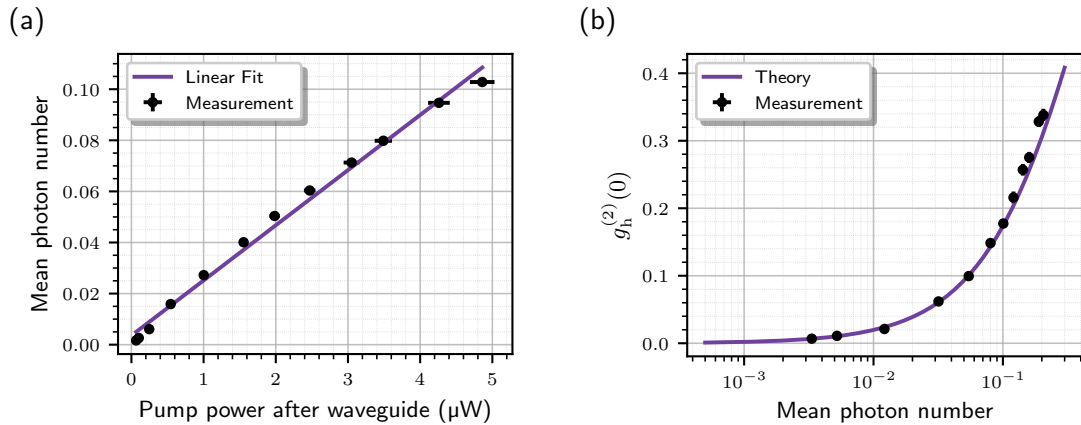


Fig. 5.10: (a) Mean photon number as a function of the pump power with an applied linear fit. (b) Heralded second-order correlation function $g_h^{(2)}(0)$ as a function of the mean photon number. The theoretical curve is in excellent agreement with the measurement.

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so, we have to take three factors into account which are the reflectivity at the end face of the waveguide, the reflectivity at the lens as well as the coupling efficiency. To estimate the reflectivity at the end face of the waveguide, we use the standard reflectivity of LN of 14.54 % at 534 nm. The fraction of light reflected at the lens is estimated via the refractive index of the lens material. For this, we consider a reflectivity of 5 %. Lastly, we have to estimate the coupling efficiency. Here, we simulate the overlap between the waveguide mode and the mode distribution given by the lens. This yields a coupling efficiency of 40 %. Summing up all of this factors leads to an overall transmission from the waveguide to the power meter of 32 %. This value is used as a correction factor to estimate the pump power inside the waveguide from the measured pump power after the waveguide. Thus, a pump power of 544 nW after the waveguide equals a pump power inside the waveguide of $P = 1.68 \mu\text{W}$. To estimate the brightness, we need next to the pump power the PDC rate and the bandwidth of the generated photons. We adjust the PDC rate with the Klyshko efficiency to get the number of generated pairs inside the waveguide which leads to a PDC rate of $R_{\text{PDC}} = 15.89 \text{ MHz}$. Lastly, we determine the bandwidth of the generated signal photons from the measured spectrum to be $\Delta\omega = 2104 \text{ GHz}$. These values lead to a brightness of $0.44 \cdot 10^7 \frac{\text{pairs}}{\text{s} \cdot \text{mW} \cdot \text{GHz}}$.

Comparing this value to other reported PDC sources, we can observe that our source features a brightness that is three order of magnitude larger than that of bulk crystal sources [157] as well as third-order nonlinear material sources such as silicon or silicon nitride [158]. This increase in the brightness is due to the longer interaction length in waveguides for the bulk crystal source and due to the higher second-order nonlinearity compared to a third-order nonlinearity. Moreover, the TFLN PDC source is two orders of magnitude brighter than a PDC source realized in weakly guiding waveguides [159]. This can be attributed to the increased mode overlap between the interacting modes in contrast to the overlap in weakly guiding waveguides. In comparison to other single pass PDC sources realized in TFLN [138, 134], we achieve a similar brightness but with the feature of generating photons at two very distinct wavelengths.

To summarize, in this section we discussed the design and experimental demonstration of a highly non-degenerate PDC source in TFLN. We show how to design a sample and include possible fabrication imperfections. Furthermore, we analyze the spectral characteristics and we could indicate an excellent agreement between our simulation and the experiment. Moreover, we evaluate the photon number purity and demonstrated that the source can generate high quality heralded single photons by analyzing the heralded second-order correlation function. Lastly, we analyzed the brightness of our source which is comparable to other PDC source in TFLN. This source is a fundamental building block for a variety of applications in the field of integrated quantum optics.

5.3 Thin-Film Thickness Compensation

In the previous section, we have already seen that the waveguide geometry has an immense impact on the dispersion. This means for example for a nonlinear process that a constant change in one geometry parameter leads to a shift of the generated output wavelength for a fixed poling period. Among all parameters, the thin-film thickness has the highest impact on the dispersion [133, 150] (see Figure 5.3). To understand the impact of the thin-film thickness on the output wavelength, we investigate a type 0 PDC process pumped with a

5.3. Thin-Film Thickness Compensation

CW-laser at 775 nm to generate both photons at 1550 nm. We choose the initial geometry to be: $t=605$ nm, $h=300$ nm, $w=1$ μ m, $a=30^\circ$, $c=0$ μ m. For this geometry and process, we simulate the JSI and plot a projection onto the signal axis in Figure 5.11. Here, we calculate the poling period with the initial geometry and simulate the output spectrum for a thinner ($t = 600$ nm) and thicker ($t = 610$ nm) thin-film thickness. A change of only 5 nm already leads to a wavelength shift of more than 50 nm. This shows that a constant change in the geometry leads to an offset in the generated output wavelength.

However, it is also possible that we do not have a constant change in one geometry parameter, but that the geometry parameter changes along the waveguide. This can be observed when measuring the thin-film thickness of a 1 cm x 1 cm sample with an ellipsometer (see Figure 5.12). We can see that the thin-film thickness changes from 597 nm to 610 nm. Moreover, some regions of the sample show less fluctuations (upper part) than other regions (lower part). Thus, depending on the position of the waveguide, the phase matching function is more or less influenced by the variations.

To simulate the influence of a varying thin-film thickness on the PDC spectrum, we cannot assume anymore that $\Delta\beta$ in Equation 2.114 is independent of the position. Now, $\Delta\beta$ is a function of the position and we have to use a generalized expression [63]:

$$A_3(L) \propto e^{i \int_0^L \Delta\beta(\zeta) d\zeta}. \quad (5.7)$$

We solve this equation via a numerical approach and use the Riemann sum as the simplest method for integration [160, 148]:

$$\phi \propto \frac{1}{L} \int_0^L e^{i \int_0^z \Delta\beta(\zeta) d\zeta} dz \approx \frac{\Delta z}{L} \sum_{m=0}^N e^{i \Delta z \sum_{m=0}^n (\Delta\beta_m - \frac{2\pi}{\Lambda})}. \quad (5.8)$$

This leads to the fact that the phase matching function can vary drastically in its shape and therefore does not necessarily exhibit a sinc² profile anymore.

To investigate the impact of a varying thin-film thickness on the shape of the phase matching function, we simulate the projection onto the signal axis of the JSI including a varying

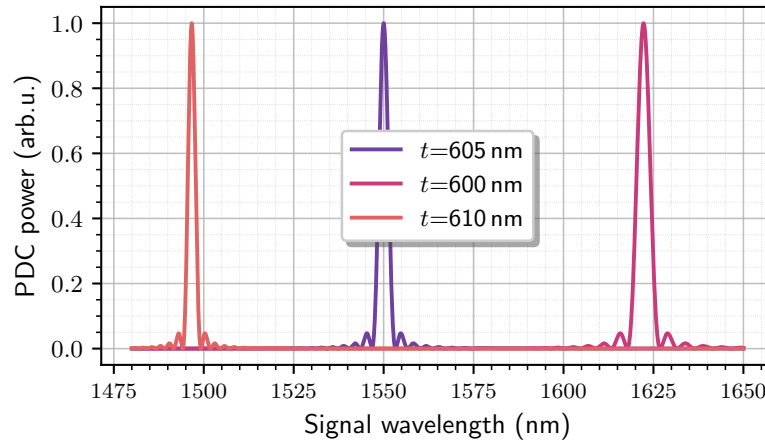


Fig. 5.11: PDC power as a function of the signal wavelength for three different thin-film thicknesses t , simulated for a fixed poling period corresponding to a thin-film thickness of 605 nm. A change of only 5 nm in the thin-film thickness has a significant impact on the generated signal and thus also on the idler wavelength.

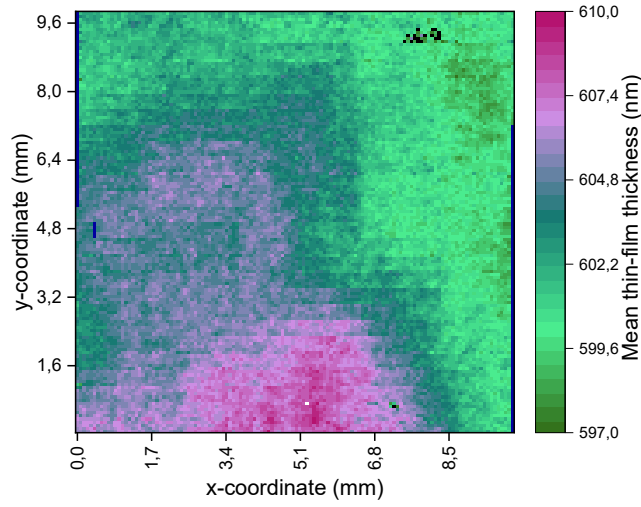


Fig. 5.12: Thin-film thickness of a 1 cm $\times$ 1 cm TFLN sample measured with an ellipsometer. The thin-film thickness varies from 597 nm to 610 nm. The black spots are measurement artifacts.

effective refractive index due to the thin-film thickness variations via Equation 5.8. The simulation code has been developed during the master thesis of Ernst-Lukas Kuhlmann [161]. To do so, we take a line scan from Figure 5.12 and plot the thin-film thickness as a function of the length of the waveguide (see Figure 5.13 (a)). Here, we assume that the length of the waveguide is the same as the length of the poled area. We calculate the PDC spectrum for a constant thin-film thickness using the mean value of the measured thin-film thickness and we calculate the PDC spectrum for the measured thin-film thickness by taking the changing effective refractive index into account. For a varying thin-film thickness, the PDC spectrum does not resemble a sinc^2 -shape anymore (see Figure 5.13 (b)). Moreover, also the height of the maximum decreased significantly compared to the spectrum from a constant thin-film thickness. Due to the varying thin-film thickness, the propagation constant changes and thus the poling period supports the efficient generation of light at different wavelength combinations and this leads to a distorted phase matching.

Moreover, the impact of a varying thin-film thickness depends on the length of the poled area and the kind of thin-film variation. To investigate this behavior, we simulate the evolution of the PDC spectrum for different length of the waveguide from 2 mm to 10 mm in 2 mm steps in Figure 5.14. The small variations within the first 2 mm do not have a significant impact on the spectrum (see Figure 5.14 (a)). This is due to the fact that the overall thin-film thickness is constant and we only have small fluctuations around a slowly varying mean value. This variation in the thin-film thickness resembles white noise and we can conclude that such kind of noise does not have a significant impact. More information about the influence of white noise and $1/f$ noise variations in the thin-film thickness as well as the top width and etching depth can be found in [149]. The situation changes when we look at a 4 mm waveguide. The first 3 mm of the thin-film have again only small variations similar to white noise as before. However, from 3 mm onward the thin-film thickness increases from 608 nm to 610 nm. This significant change in the thin-film thickness leads to a reduction

5.3. Thin-Film Thickness Compensation

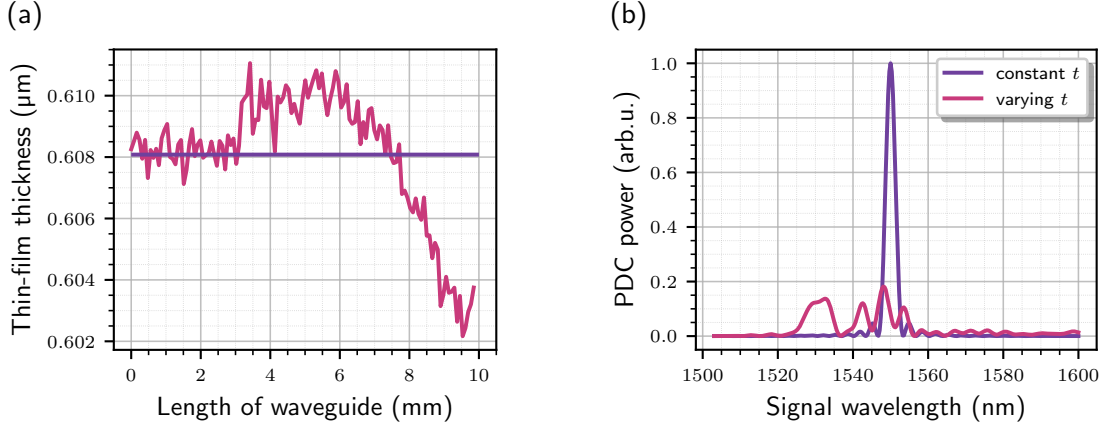


Fig. 5.13: (a) Measured thin-film thickness as a function of the waveguide length. The horizontal line indicates the mean thin-film thickness. (b) PDC spectrum as a function of the signal wavelength for a constant thin-film thickness and for a varying thin-film thickness. It can be seen that a varying thin-film thickness leads to a destruction of the sinc^2 -shape.

of the main peak to below 70 % due to the fact that the varying thin-film thickness leads to different dispersion resulting in different phase matching wavelengths for a fixed poling period, see Figure 5.14 (b). Moreover, we can see that we get more pronounced side peaks at shorter wavelengths. An even more drastic difference between the phase matching function of a constant thin-film and a varying thin-film can be observed in Figure 5.14 (c) which resembles a 6 mm long waveguide. Here, we see two peaks with an equal height of around 30 %. This is due the fact that we have two main thin-film thicknesses namely 608 nm and 610 nm which lead to two phase matched wavelengths. Using even longer waveguides leads to an even more distorted spectrum (see Figure 5.14 (d) and (e)), since the variations in the thin-film thickness resemble an almost linear decrease from 610 nm to 602 nm leading to numerous phase matching wavelengths for a fixed poling period. This shows that the shape of the PDC spectrum strongly depends on the position of the waveguide on the sample (placed in an area with more or less variations), on the kind of thin-film variation and on the length of the waveguide/the length of the poled area. Such a PDC spectrum as shown in Figure 5.14 (e) is unwanted since it destroys the shape as well as reduces the efficiency of the process. Different methods have been proposed and demonstrated to counteract the change in the effective refractive index due to the thin-film variations which will be discussed next.

5.3.1 State of the Art - Thin-Film Thickness Compensation

To counteract the thin-film thickness variations in a nonlinear process, we need to adapt the momentum conservation condition:

$$\Delta\beta = \beta_1 + \beta_2 + \beta_3 + \frac{2\pi}{\Lambda} \approx 0. \quad (5.9)$$

This means, we can either change the propagation constants or we can change the poling period.

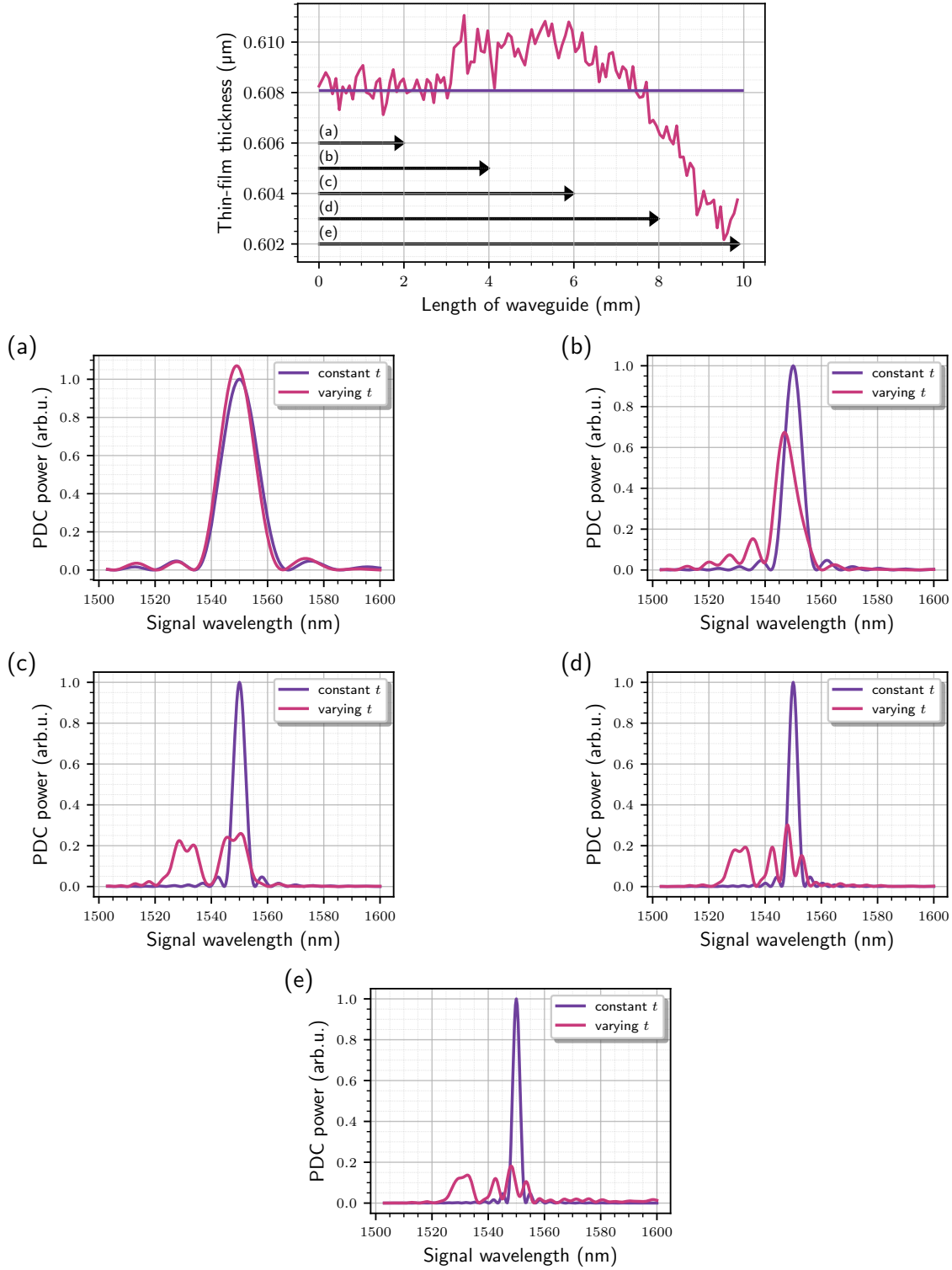


Fig. 5.14: Evolution of the PDC spectrum for different length of the poling. The upper picture shows a measured thin-film thickness. The plots (a)-(e) show the resulting signal spectrum for 2 up to 10 mm in 2 mm steps. With increasing length of the waveguide, the shape of the spectrum resembles less and less a sinc^2 shape and the main peak power reduces.

5.3. Thin-Film Thickness Compensation

Changing the poling period to counteract thin-film thickness variations was demonstrated by Chen et al. in 2023 [150]. In this method, the poling period is calculated locally for each thin-film thickness. This means, if the thin-film thickness changes, the effective refractive index of the waveguide is altered due to the high confinement which leads to a different poling period for a fixed wavelength combination. Therefore, the poling period is locally adjusted to the thin-film thickness to achieve the generation of photons at the same wavelengths.

Instead of changing the poling period locally which is challenging to fabricate, we can also alter the effective refractive index via the temperature. This has been demonstrated by Li et al. in 2024 [162] using the fact that the refractive index is temperature dependent. To change the effective refractive index of the waveguide, they used segmented micro-heaters. By adjusting the temperature to the LN thin-film locally, the effective refractive index is altered such that it matches again to the poling period. However, to realize this an extra fabrication step is needed. To avoid this, we can also change the effective refractive index by altering the top width during the fabrication.

In this method, we change the top width of the waveguide in relation to the thin-film thickness. The simultaneous change results in a constant effective refractive index and thus, a constant poling period can be applied. We filed a patent about this method [163]. Shortly after, He et al. theoretically demonstrated this method in a publication [164]. In the next section, we will elaborate on this method.

5.3.2 Adapted Width

The top width is an ideal parameter to counteract thin-film thickness variations. We can

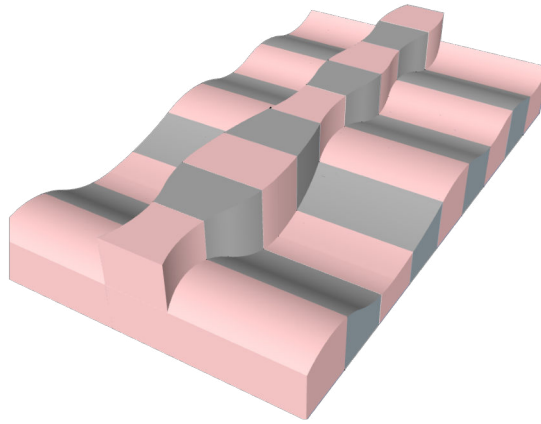


Fig. 5.15: Conceptual drawing of the adaptive width approach. The image shows a periodically poled waveguide with a varying thin-film thickness. To counteract the variations in the thin-film thickness, the top width of the waveguide is changed accordingly.

measure the thin-film thickness before the waveguide fabrication along a line and simulate a corresponding top width for each thin-film thickness to match the effective refractive index along the waveguide. Afterwards, we can fabricate a waveguide with tapers to change the top width locally along the waveguide. This is shown in Figure 5.15 in which we show a periodically poled waveguide with a varying thin-film thickness and an adjusted top width. To investigate the relation between the top width and the thin-film thickness, we start

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with simulations. We calculate the phase mismatch $\Delta\beta_0$ for a type 0 PDC process from 775 nm to 1550 nm for the following initial geometry: $t = 608$ nm (mean value of measured thin-film thickness in Figure 5.13 (a)), $w = 1 \mu\text{m}$, $h = 300$ nm, $a = 30^\circ$, $c = 0 \mu\text{m}$. Based on this geometry, we calculate the phase mismatch $\Delta\beta$ for a range of varying top widths and thin-film thicknesses. Finally, we plot the phase mismatch difference $\Delta\beta - \Delta\beta_0$ as a function of the top width and the thin-film thickness as shown in Figure 5.16. If the phase mismatch difference is zero, the geometries have the same effective refractive index. In the plot, we can observe that an increasing thin-film thickness needs a decreased top width to preserve the effective refractive index. This confirms that we can use the top width to counteract thin-film thickness variations.

To do so, we have to find a corresponding top width for each thin-film thickness. In Figure 5.17 (a) we show the thin-film thickness measured with an ellipsometer along a line as well as the top width along the waveguide which is needed to counteract the thin-film variations. As seen above, a decreasing thin-film thickness can be compensated by an increased top width. To counteract a change of around 8 nm in the thin-film thickness, a change in the top width of 50 nm is needed. If we simulate the signal spectrum of a constant thin-film thickness and a varying thin-film thickness compensated via an adaptive top width, we can reconstruct the sinc^2 -shape as well as the peak height (see Figure 5.17 (b)). Thus, adapting the thin-film thickness via the top width shows strong evidence to be a promising method which needs to be demonstrated experimentally.

This concludes the section about first-order dispersion engineering to realize source with a desired wavelength combination as well as a brief discuss about the influence of thin-film thickness variations. In the last section of this chapter, we want to investigate theoretically PDC sources based on second-order dispersion engineering.

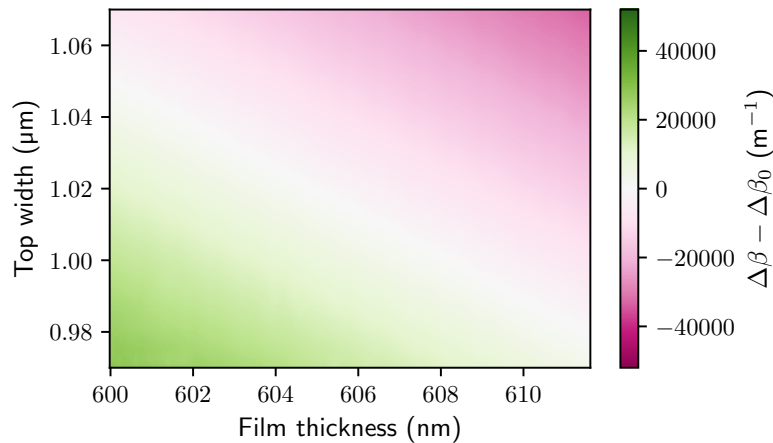


Fig. 5.16: Relation between top width and thin-film thickness as a function of the phase mismatch $\Delta\beta - \Delta\beta_0$. We can see that for each thin-film thickness we can find a top width such that the phase mismatch is again zero.

5.4. Second-Order Dispersion Engineering

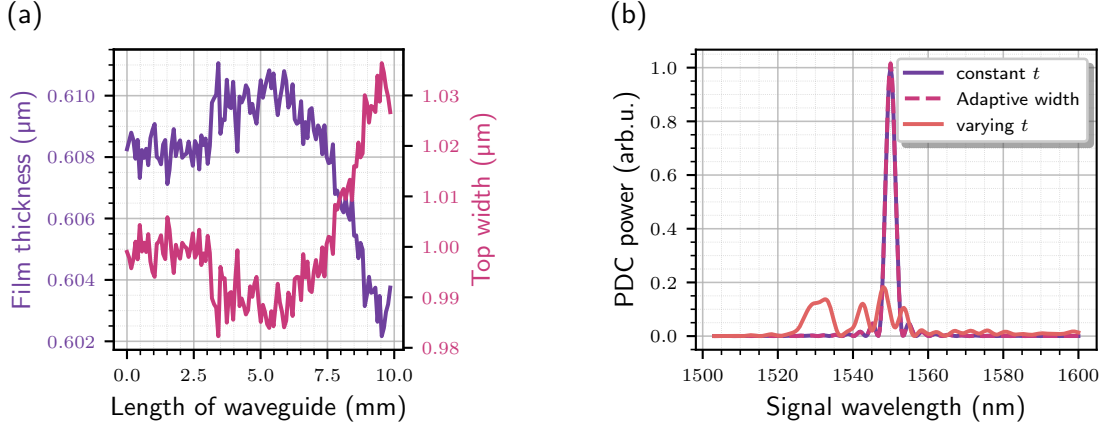


Fig. 5.17: (a) Measured thin-film thickness as a function of the length of the waveguide. Moreover, we also show the top width as a function of the length of the waveguide necessary to counteract the thin-film variations. (b) PDC spectrum with a constant thin-film thickness as well as with a varying thin-film thickness combined with the adaptive top width approach and with a varying thin-film thickness without any compensation. We can see that we can reconstruct the sinc^2 shape by adapting the top width of the waveguide to the thin-film thickness variations.

5.4 Second-Order Dispersion Engineering

In this section, we will discuss second-order dispersion. This means we will not only investigate the phase velocities as we did with first-order dispersion engineering to achieve a PDC source with a desired wavelength combination, but we will also tailor the group velocities to achieve phase matching functions with specific phase matching angles. Here, we will focus on type 2 PDC sources pumped at 775 nm to generate the signal and idler photon at 1550 nm. These types of sources are of great interest since the generated photons at this wavelength can be transmitted via the existing fiber network with low losses. Moreover, with this configuration a photon pair source without correlations between the signal and idler photons can be achieved which is crucial for any quantum interference experiments and is a requirement for quantum computing. Two particular PDC sources are crucial as already discussed in section 2.5.1. The first one is a PDC source with asymmetric group velocity matching (0° phase matching angle) and the second one is a symmetric group velocity matched source (45° phase matching angle). Both sources have various applications in the field of (integrated) quantum optics. An asymmetric group velocity matched source [73] can be used to generate heralded pure single photons as well as to herald pure single photons in arbitrary temporal shapes [165]. A symmetric group velocity matched source [126] can generate two indistinguishable photons which are useful for any type of quantum interference and are prerequisite for quantum computing such as Gaussian boson sampling [166].

Thus, here we want to focus on these two kinds of sources. For this, we will first discuss the state of the art of dispersion engineering single pass PDC sources pumped at 775 nm to generate photons at 1550 nm in TFLN waveguides.

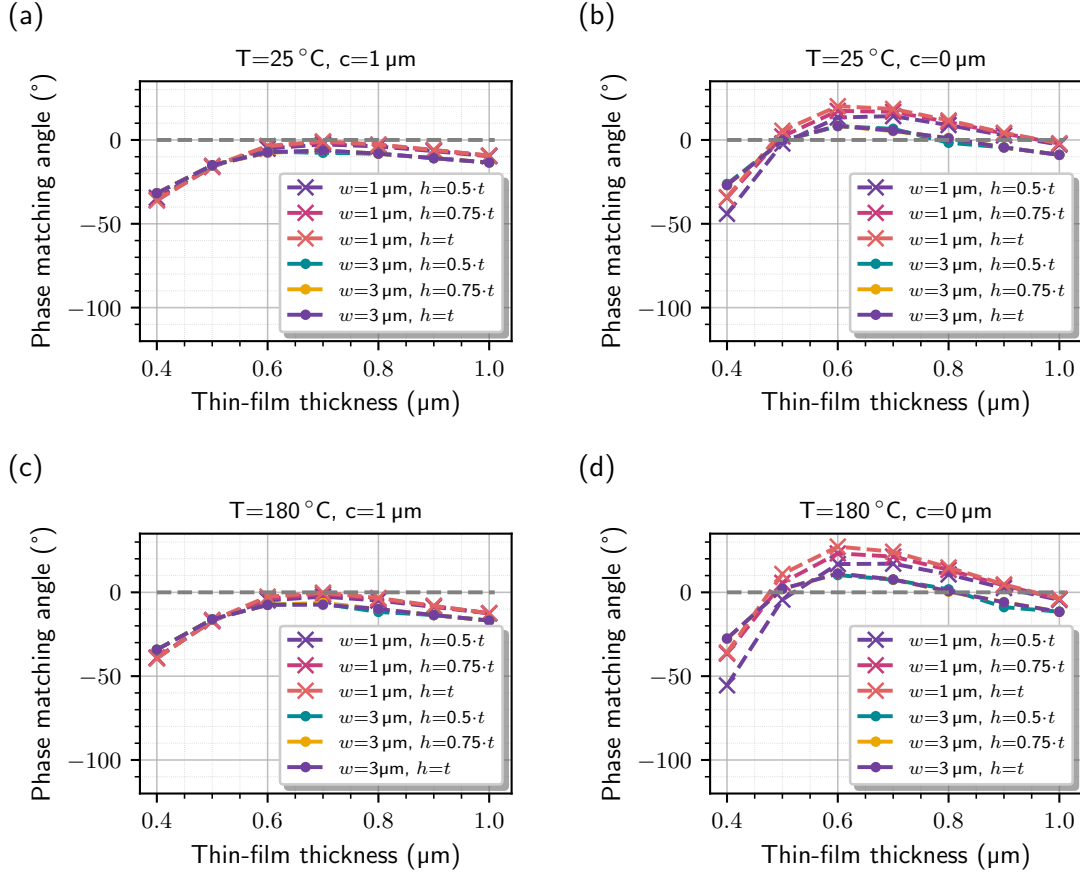


Fig. 5.18: Phase matching angle of a type 2 PDC process pumped at 775 nm to generate both photons at 1550 nm. (a) With a cladding layer at 25°C. (b) Without a cladding layer at 25°C. (c) With a cladding layer at 180°C. (d) Without a cladding layer at 180°C.

5.4.1 State of the Art - Second Order Dispersion Engineering

The first dispersion engineered type 2 PDC source in TFLN was demonstrated by Xin et al. [137]. They realized a PDC source with an almost horizontal phase matching function realizing a close to asymmetric group velocity matched PDC source in TFLN. Another PDC source with an almost horizontal phase matching function was demonstrated by Kumar et al. also in 2022 [167]. The authors measured a phase matching angle of -3° , but they did not show any experimental quantum measurements. Thus, single-pass asymmetric group velocity matched PDC sources in TFLN have already been investigated. However, a PDC source with symmetric group velocity matching has not been shown yet. Here, we will propose two geometries to achieve such a type of PDC source in TFLN.

5.4.2 Symmetric Group Velocity Matched PDC Source in TFLN

To find a suitable waveguide geometry for a symmetric group velocity matched PDC source in TFLN, a broad parameter sweep is performed. We investigate different thin-film thicknesses

5.4. Second-Order Dispersion Engineering

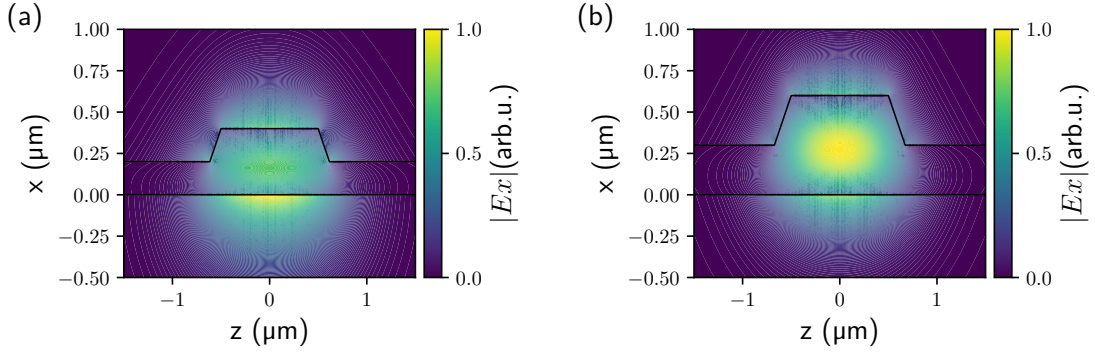


Fig. 5.19: Simulated mode distribution of a cross-section of a waveguide for a thin-film thickness of (a) 400 nm and an etching depth of 200 nm and (b) for a thin-film thickness of 600 nm and an etching depth 300 nm at 1550 nm. Both waveguides have a top width of 1 μm and a waveguide angle of 30° . The mode in the thinner film is not as well guided as the mode in the thicker film.

from 400 nm to 1000 nm in 100 nm steps. For each thin-film thickness, we simulate different etching depths, namely 50 %, 75 % and 100 % of the thin-film thickness, for a 1 μm and for a 3 μm top width. The waveguide angle is fixed to 30° . We simulate these geometries with a 1 μm SiO_2 cladding layer and without a cladding layer at room temperature and at 180°C . For all of these geometries, we calculate the phase matching angle via Equation 2.125 for a type 2 PDC process pumped at 775 nm to generate the photons at 1550 nm. We plot the phase matching angle as a function of the thin-film thickness in Figure 5.18. In all plots, we show the waveguide geometries with a top width of 1 μm marked as dots and with a top width of 3 μm as crosses. We connect the individual points with a line to guide the eye. In Figure 5.18 (a) we analyze geometries at room temperature with a 1 μm thick cladding layer. The phase matching angle for a thin-film thickness of 400 nm for all simulated geometries is between -40° and -30° . With increasing thin-film thickness, the phase matching angle increases to slightly below 0° at a thin-film thickness of 700 nm. For more than 700 nm thin-film thickness, the phase matching angle decreases to around -10° . The negative phase matching angle for thinner thin-film thicknesses is due to the fact that the TM modes at 1550 nm are not completely guided in the lithium niobate thin-film compared to modes for example in a 600 nm thick thin-film (see Figure 5.19). Increasing the thin-film thickness leads to a better confinement of the TM mode which results in an increase of the phase matching angle. At thicker thin-film thickness the phase matching angle decreases again, since the modes are less confined. Thus, the dispersion is less influenced by the geometry and more by the material dispersion. Removing the cladding layer, increases the phase matching angle (see Figure 5.18 (b)) due to the fact that the index contrast is increased. Already for a thin-film thickness of 500 nm, the phase matching angle is above 0° . The maximal phase matching angle of around 20° is reached at a thin-film thickness of 600 nm. For thicker thin-films the phase matching angle starts to reduce again. Moreover, we can again observe that the etching depth has only a significant influence on the 1 μm top width waveguides and not on the 3 μm top width waveguides, as already seen in Figure 5.1. From the plot, we can conclude that to achieve a high phase matching angle a small top width

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and a deep etching depth is necessary. However, with this configuration we cannot reach symmetric group velocity matching. Another parameter that can alter the dispersion is the temperature. In Figure 5.18 (c) and (d) we plot the phase matching angle as a function of the thin-film thickness at 180 °C with and without a cladding layer, respectively. We can see that the phase matching angle increase to around 25 ° for a thin-film thickness of 600 nm without a cladding layer. This is due to the fact that the refractive index of lithium niobate increases with increasing temperature leading to a higher index contrast change. Applying a cladding layer reduces again the phase matching angle due to the lower index contrast. This study leads to the conclusion that we might achieve a symmetric group velocity matched PDC source in TFLN at a high temperature with a narrow top width and a deep etching depth at 600 nm thin-film thickness. Thus, we will next investigate the temperature dependency of such geometries in detail.

Tuning the Temperature

We have already seen that the phase matching angle is temperature dependent and that a narrow top width together with a deep etching depth achieve a high phase matching angle. Thus, we simulate waveguides with a narrower top width of 800 nm and again 1 μm for different temperature between 25 °C and 300 °C for two different etching depths (450 nm and 600 nm). In Figure 5.20 the phase matching angle as a function of the temperature for each geometry is displayed. The plot shows that an increase in the temperature leads to an overall increase in the phase matching angle. With a top width of 1 μm , we can change the phase matching angle from around 20 ° to above 30 ° if we increase the temperature to 300 °C. For the smaller waveguide top width of 800 nm, the phase matching angle increases even more to 40 ° for a etching depth of 450 nm and to 50 ° for 600 nm etching depth. Thus, with a small waveguide top width and a completely etched thin-film a symmetric group velocity matched PDC source is reachable at high temperatures. One of those geometries that provides a 45 ° phase matching angle is the following: thin-film thickness $t=600$ nm, etching depth $h=600$ nm, top width $w=800$ nm, angle $a=30$ °, cladding thickness $c=0$ μm

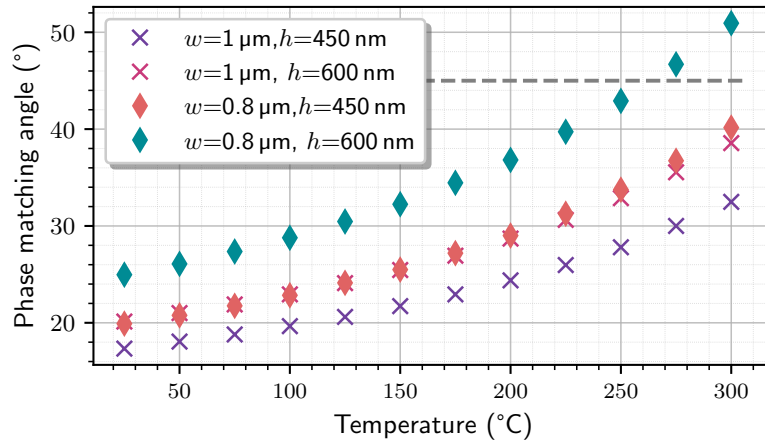


Fig. 5.20: Phase matching angle as a function of the temperature for different waveguide geometries. The higher the temperature, the higher the phase matching angle. We plot a horizontal line at 45° phase matching angle which can be reached for a 800 nm top width and a completely etched waveguide at above 275 °C.

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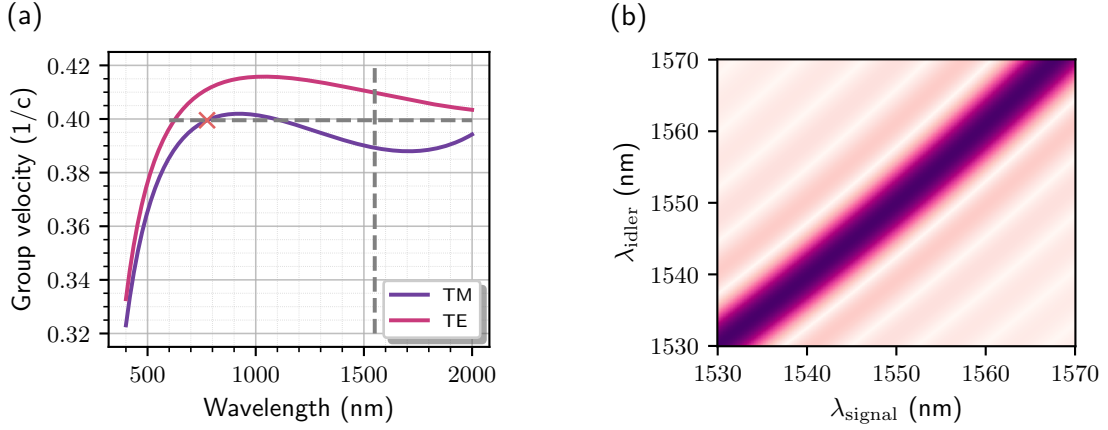


Fig. 5.21: Dispersion of a waveguide geometry at high temperatures which yields a phase matching angle of 45° C. (a) Group velocity as a function of the wavelength. We marked the group velocity of the pump with a mark. The vertical line is at the signal and idler wavelength and crosses a vertical line at the height of the group velocity of the pump photon. (b) Phase matching function in the signal and idler plane for a 5 mm poling length.

and temperature $T=275^\circ\text{C}$.

In Figure 5.21 (a) we show the group velocity as a function of the wavelength for this geometry. The horizontal line indicates the group velocity at the pump wavelength and the vertical line is placed at 1550 nm. Here, we can see that the group velocity of the pump wavelength is in the middle of the group velocity of the signal and idler wavelength. This symmetry leads to a 45° phase matching angle which we show in Figure 5.21 (b) where we used a poling length of 5 mm.

However, this geometry requires on the one hand a small waveguide top width leading to high losses as discussed in chapter 3 and on the other hand needs a high temperatures which is not always ideal, especially for real life applications. Therefore, we looked at a different approach for realizing a 45° phase matching angle such as freestanding or suspended waveguides.

Suspended Waveguides

Waveguides which are not attached to the SiO_2 layer underneath the lithium niobate film are called free-standing waveguides or suspended waveguides [168, 169, 170]. Those waveguides are already used in TFLN for Brillouin scattering, SHG and electro-optic modulators [171, 172, 173, 168].

Here, we use these waveguides to alter the dispersion. Removing the SiO_2 underneath the lithium niobate film completely or partially (see Figure 5.22 (a)) leads to an increased index contrast at the bottom of the waveguide. Now, the index contrast is between lithium niobate and air instead of lithium niobate and SiO_2 . To investigate the dispersion of such a waveguide structure in more detail, we simulate waveguides geometries with a 600 nm thin-film thickness and an angle of 30° for different etching depths and top widths at room temperature. For all these geometries, we calculate the phase matching angle via Equation 2.125 for a type 2 PDC process pumped at 775 nm to generate both photons at 1550 nm. In Figure 5.22 (b), the phase matching angle as a function of the top width

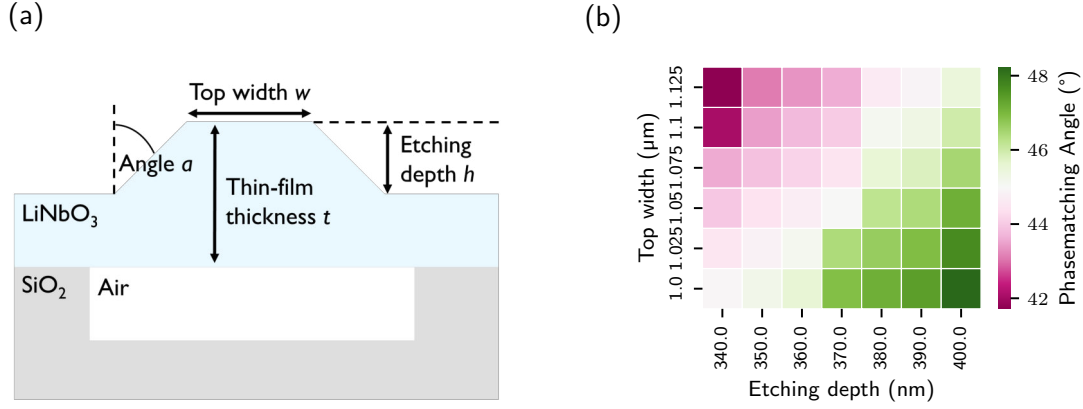


Fig. 5.22: (a) Sketch of a cross-section of a suspended waveguide in TFLN. (b) Phase matching angle has a function of the top width and etching depth for suspended waveguides. A phase matching angle of 45° is reachable.

and etching depth is shown. For a top width between $1\ \mu\text{m}$ and $1.125\ \mu\text{m}$ and an etching depth between $340\ \text{nm}$ and $400\ \text{nm}$, the phase matching angle is between 42° and 48° . This means with a suspended waveguide we can achieve symmetric group velocity matching in a TFLN waveguide. One geometry that leads to a phase matching angle of 45° is the following: thin-film thickness $t=600\ \text{nm}$, etching depth $h=370\ \text{nm}$, top width $w=1.05\ \text{nm}$, angle $a=30^\circ$, cladding thickness $c=0\ \mu\text{m}$, temperature $T=25^\circ\text{C}$.

The group velocities of this geometry as a function of the wavelength are shown in Figure 5.23 (a). Again, the horizontal line indicates the group velocity of the pump wavelength and the vertical line is at $1550\ \text{nm}$ showing the symmetric arrangement of the TE and TM group velocity at $1550\ \text{nm}$ in comparison to the pump wavelength. In Figure 5.23 (b) the phase matching function in the signal and idler plane is shown with a poling length of $5\ \text{mm}$. Comparing this phase matching function to the one in Figure 5.21 (b), we can observe that the phase matching bandwidth is smaller. The phase matching bandwidth for the normal waveguide at higher temperatures is $3.14\ \text{nm}$. For the suspended waveguides, the phase matching bandwidth is only $1.37\ \text{nm}$. This reduced bandwidth is due to the difference in the group velocities at the pump wavelength and the signal and idler wavelength. The waveguide at high temperature has a difference between the pump and signal/idler group velocity of around $3 \cdot 10^6\ \text{m/s}$. The difference between the pump and signal/idler group velocity is for the suspended waveguides $7 \cdot 10^6\ \text{m/s}$ which is more than twice as large. This greater difference in the group velocity leads to a narrower bandwidth of the phase matching function for the same poling length.

To sum up, we can achieve a 45° phase matching angle in TFLN waveguides by either heating the sample or using suspended waveguides. We filed a patent about a PDC source in TFLN based on suspended waveguides [174].

This chapter shows that TFLN is an ideal material to shape the dispersion via the geometry. By changing the etching depth of the waveguide, we can tune the effective refractive index from 1.909 to 1.876 at $1550\ \text{nm}$ due to the high mode confinement. This change in the effective refractive index leads consequently to a change in the group velocities. Due to the broad tuning range in TFLN waveguides, we can alter the slope of the group velocities from

5.4. Second-Order Dispersion Engineering

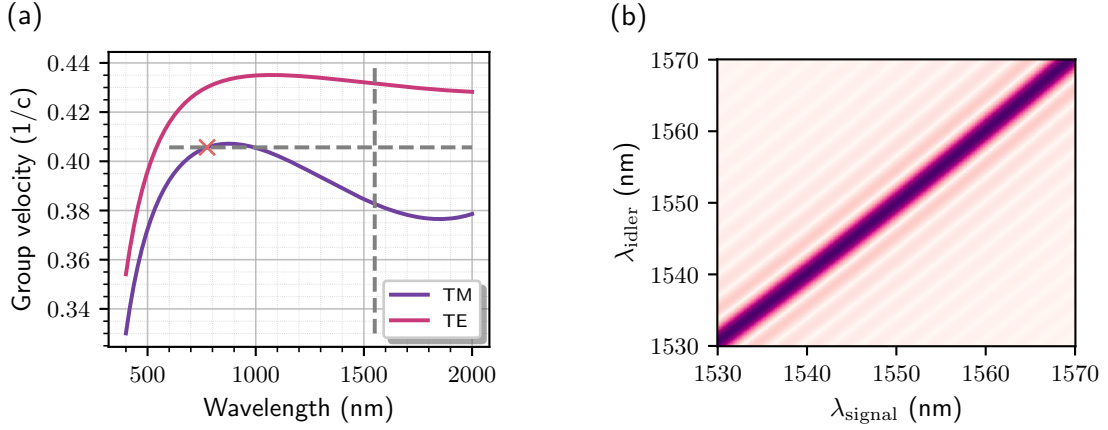


Fig. 5.23: Dispersion of a suspended waveguide geometry which yields a phase matching angle of 45° . (a) Group velocity as a function of the wavelength. A horizontal line at the height of the group velocity of the pump wavelength (marked with a cross) shows together with the vertical line at 1550 nm, the symmetric group velocities of the signal and idler wavelength around the pump wavelength. (b) Phase matching function in the signal and idler plane for a suspended waveguides for a 5 mm poling length.

a negative to a positive slope. This immense variation leads to the possibility of achieving PDC source with a phase matching angle between -40° and more than 45° meaning we can realize a asymmetric and a symmetric group velocity matched PDC source. Thus, the dispersion engineering capabilities of TFLN waveguides open up the possibility of achieving different types of PDC sources for various applications.

The here experimentally and theoretically demonstrated source are important building blocks for integrated quantum optics circuits. The non-degenerate PDC source can be used as a heralded single photon source, but is also important for quantum spectroscopy applications. The theoretically investigated waveguide geometries for a symmetric group velocity matched PDC source are highly relevant for any kind of quantum interference experiment. Such sources are required for quantum computing such as Gaussian boson sampling. Thus, the experimental demonstration of such a source in TFLN paves the way for a integrated quantum processor on chip. Therefore, we can revolutionize the field of integrated quantum optics by realizing dispersion engineered PDC sources in TFLN waveguides.

Summary

In this chapter, we discussed dispersion engineering of PDC sources in TFLN. We started with first order dispersion engineering in which we matched phase velocities to achieve a PDC source with a desired wavelength combination. Here, we showed the design and the detailed experimental characterization of a highly non-degenerate PDC source in TFLN. During the design of the PDC source, we discussed the influence of fabrication tolerances which we continued in more detail by discussing the influence of thin-film thickness variations on the phase matching function. The impact of the thin-film thickness on the PDC spectrum was investigated and we introduced methods to counteract the thin-film thickness variations.

Chapter 5. Engineering TFLN Photon Sources

Lastly, we showed two approaches to realize a PDC source with symmetric group velocity matching. This chapter shows that we are able to engineer and control the dispersion of TFLN waveguides in various ways.

6

Quantum Pulse Gate in TFLN

In the last chapter, we have shown the possibilities of dispersion engineering in TFLN waveguides for PDC sources. Next, we want to utilize our gained knowledge about dispersion engineering to realize the first TFLN quantum pulse gate (QPG), as introduced in subsection 2.5.3. As explained, the QPG is a device which can selectively up-convert arbitrarily chosen temporal modes to a pre-defined output frequency due to the underlying group velocity matched SFG process. Temporal modes are a versatile framework in the field of quantum information science [86, 175], since information can be encoded in the spectral degree of freedom. However, using temporal modes faces the challenge of efficient characterization and manipulation as they can overlap both in frequency and time. Since the QPG is a device which can manipulate temporal modes, it is a device useful for this research field. To date, QPGs have been mainly realized in titanium indiffused waveguides in lithium niobate. These feature a low normalized conversion efficiency limited by two restrictions. First, the low-index contrast waveguide leads to a low spatial mode overlap between the three interacting modes. Second, the material-dependent dispersion of titanium indiffused waveguides causes the realization of QPGs only for specific wavelengths and processes. The resulting lower normalized conversion efficiency hinders the application of QPGs in scalable quantum experiments, such as frequency encoded quantum networks [176]. To utilize QPGs in these applications, we need to boost the efficiency. To do so, we will exploit the high index contrast waveguides in TFLN. Furthermore, we will also harness the dispersion engineering possibility to explore the more efficient type 0 SFG process to enhance the efficiency further to demonstrate a highly efficient QPG. Thus, in this chapter, we will realize the first QPG in TFLN. For this, we will discuss the state of the art QPGs. Afterwards, we will describe how we engineer the device and explain the measurement setup. Lastly, we will benchmark the QPG in TFLN by evaluating the phase matching function, the temporal mode selectivity and the conversion efficiency. The results demonstrate that with TFLN we can boost the normalized conversion efficiency by three orders of magnitude compared to devices based on titanium indiffused waveguides in lithium niobate which paves the way for the implementation of QPGs in large scale quantum networks.

6.1 State of the Art - Quantum Pulse Gate

The first theoretical proposal of the QPG was published by Eckstein et al. in 2011 [82]. The authors introduced the basic concept of a QPG and simulated waveguide parameters to integrate the device in titanium indiffused lithium niobate waveguides. The wavelength of the input pulse was set to 1550 nm for compatibility with the existing low-loss fiber network.

Chapter 6. Quantum Pulse Gate in TFLN

Due to the material dispersion of lithium niobate, the pump pulse is at 870 nm to achieve group velocity matching between the input and pump pulse for a type 2 SFG process. Three years later, the first experimental realization of a QPG was demonstrated by Brecht et al. [81]. They implemented the QPG as proposed by Eckstein et al. and measured a temporal mode selectivity of 80 % which was only limited by the experimental setup. Moreover, they measured a conversion efficiency of 87.7 %.

Based on this realization, the QPG was used for various applications in the field of quantum information science such as quantum walks [177], quantum state tomography [178, 179, 180, 181], fundamental quantum mechanics [182], quantum metrology [87, 183], temporal mode sorting [184] and spectral bandwidth compression [185]. Moreover, Dileep et al. realized a double-stage QPG which increases the efficiency but does not work with single-photons [83]. In 2023, a novel QPG schemes was demonstrated by Serino et al. which is the multi-output QPG which can be operated with more temporal modes [186]. Moreover, combining a PDC source with a QPG leads to the possible of realizing high-dimensional temporal-mode based quantum networks [176].

All of the above mentioned devices are based on titanium indiffused waveguides. In these kinds of waveguides the dispersion is mainly determined by the material dispersion due to the low index contrast waveguides. Therefore, the QPG with the wavelength combination of around 860 nm as a pump wavelength and 1540 nm as an input wavelength can only achieve group velocity matching which is essential for a functional QPG in a type 2 SFG process. The low index contrast waveguide together with the less efficient type 2 process compared for example to a type 0 process results in a lower normalized conversion efficiency. Therefore, here, we use TFLN which offers the possibility of dispersion engineering due to the high index contrast waveguides. This offers the prospect to achieve group velocity matching for different wavelength combinations and process types to realize a highly efficient QPG. Thus, in this chapter, we will use all of the gained knowledge about engineering nonlinear devices in TFLN waveguide and show how we engineer and benchmark the first QPG in TFLN.

6.2 Dispersion Engineering of QPG in TFLN

As mentioned above, the standard QPG is based on a type 2 SFG process with a pump wavelength at 860 nm and an input wavelength in the telecom regime to achieve group velocity matching resulting in a phase matching function that is horizontal in the $(\lambda_{\text{in}}, \lambda_{\text{out}})$ -plane as introduced in subsection 2.5.3. This configuration leads to two main problems. First, to achieve the necessary group velocity matching between the input and pump pulse, we need to utilize a type 2 process which use the non-optimal d-tensor element d_{31} which is in the order of 3.7 pm/V-6.3 pm/V [187]. Second, the mode field overlap between the three interacting modes is low due to the low-index contrast waveguides. The low index contrast leads to the fact that the interacting modes differ in size as well as in the position in the waveguide [188]. The center of the pump and output mode is shifted compared to the input mode. Moreover, the mode of the pump and output is significantly smaller than the input mode. Due to the different size and position, it is only possible to achieve a small mode overlap. This results in a low normalized conversion efficiency, see Equation 2.146. We overcome these limitations by employing TFLN. Due to the small high-index contrast waveguides in TFLN, we achieve a significant higher field overlap between the interacting modes. Moreover, TFLNs dispersion engineering capabilities permit to achieve group velocity matching between the pump pulse at 860 nm and the input pulse at 1550 nm for a type 0 process which utilized the highest

6.2. Dispersion Engineering of QPG in TFLN

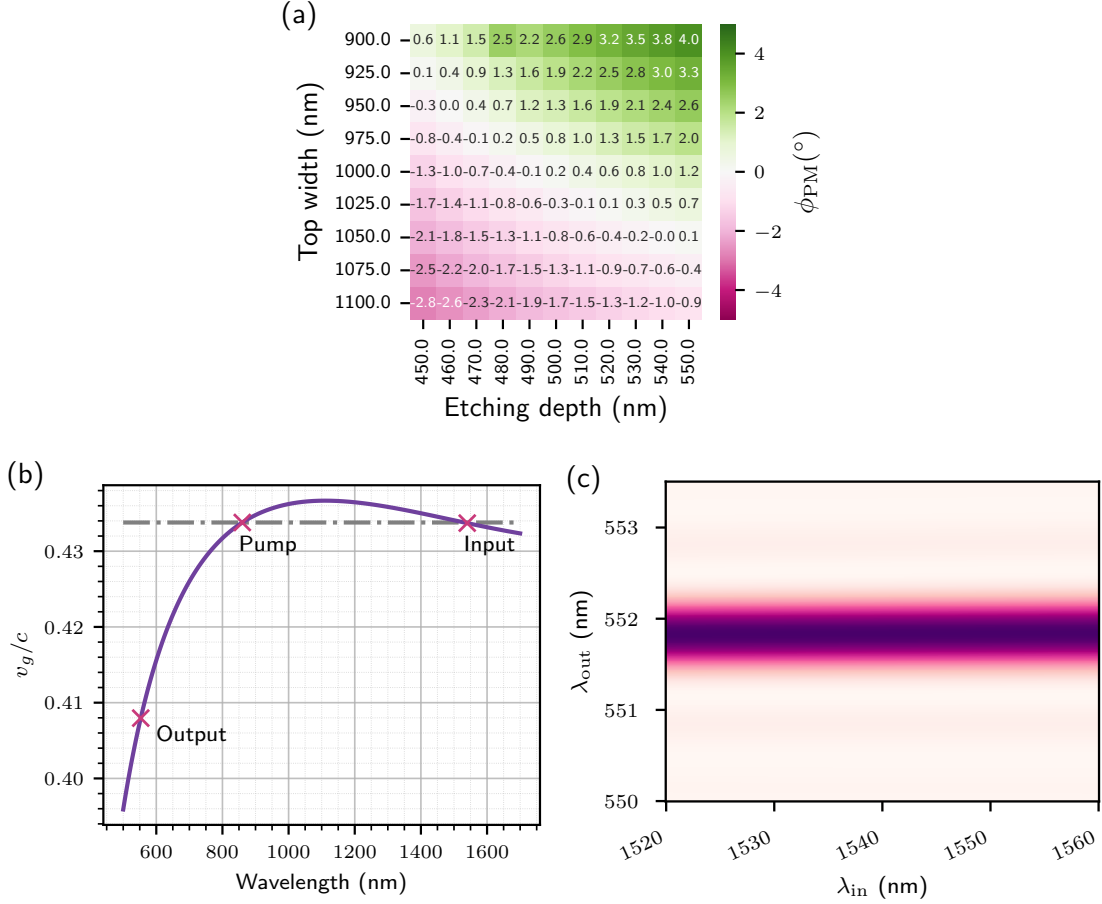


Fig. 6.1: Dispersion engineering plots for a type 0 QPG in TFLN. (a) Phase matching angle as a function of the top width and etching depth for a pump wavelength of 860 nm and an input wavelength of 1540 nm. (b) Group velocity as a function of the wavelength. The dashed grey line indicates the point of group velocity matching for the input and pump wavelength. (c) Resulting horizontal phase matching function due to the group velocity matching between the input and pump wavelength.

d-tensor element d_{33} which is in the order of 20 pm/V-47 pm/V [187]. Due to the strongly waveguide-geometry-dependent dispersion also other wavelength combinations with group velocity matching can be found. Here, we decide to use the same wavelength combination as a QPG realized in titanium indiffused waveguides for a better comparison.

To find a geometry which allows for group velocity matching between the input and pump wavelength, we simulate different top widths and etching depths for a x-cut 5mol% MgO doped TFLN sample with a thin-film thickness of 600 nm, an angle of 30° and no cladding layer. Again, we use *Lumerical* and the model explained in section 2.4.2. For each geometry, we calculate the phase matching angle ϕ_{PM} (see Equation 2.143). We plot the phase matching angle ϕ_{PM} as a function of the top width and etching depth in Figure 6.1 (a). The graph shows that in the etching depth range between 450 nm to 550 nm we can find a corresponding top width between 900 nm and 1100 nm which fulfills group velocity matching between the input and pump wavelength as indicated by a phase matching angle close to

Chapter 6. Quantum Pulse Gate in TFLN

zero. Since the losses decrease with increasing top width (compare chapter 3), we choose a top width of 1000 nm which results in an etching depth of 500 nm. In Figure 6.1 (b) the group velocity of the fundamental TE mode is plotted as a function of the wavelength. As expected the group velocity increases with increasing wavelength. But between 1000 nm and 1200 nm the group velocity starts to decrease again. This is favorable since this turning point leads to the same group velocity at 860 nm and 1540 nm as indicated by the horizontal line. Group velocity matching leads to a horizontal phase matching function as shown in Figure 6.1 (c) where we plot the phase matching function as a function of the input and output wavelength for a 3 mm long poled area. Thus, this waveguide geometry ($t=600$ nm, $h=500$ nm, $w=1$ μ m, $a=30^\circ$, $c=0$ μ m) exhibits the necessary dispersion to realize a type 0 QPG in TFLN.

Based on the simulations, we design a periodically poled sample as discussed in subsection 5.2.2 to incorporate fabrication tolerances. According to this, we choose the following poling periods for the QPG sample: 3.008 μ m, 3.093 μ m, 3.178 μ m, 3.274 μ m and 3.369 μ m. To characterize and benchmark the sample, we built a new setup which is explained in the next section.

6.3 Experimental Setup

The QPG acts on temporal modes, thus, we need to be able to shape the input and pump mode to check the mode selective behavior of the device. To do this, we build the setup shown

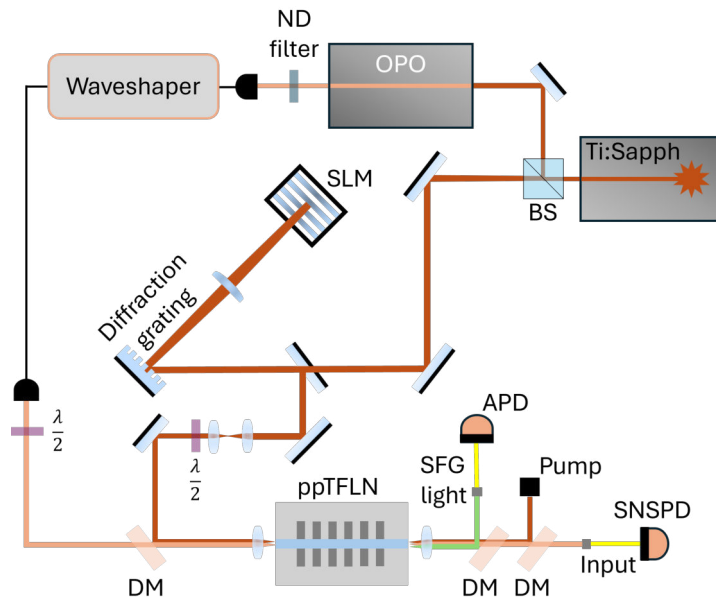


Fig. 6.2: Setup to characterize a QPG. The heart of the setup are two waveshapers, one commercial for the input light and one home-built one for the pump light, which are used to shape both pulses in different mode shapes. The shaped pulses are coupled to the periodically poled TFLN sample. The generated SFG light is measured with an APD or a single-photon sensitive spectrometer and the transmitted input light can be measured with an SNSPD.

6.4. Benchmarking of a Quantum Pulse Gate

Figure 6.2 with two pulseshapers as well as the pulsed laser system as the key components. It starts with a titanium-sapphire oscillator (Ti:Sapph) which generates ultrashort laser pulses with a repetition rate of 80 MHz at a central wavelength of $\lambda_p=872.57$ nm. After the Ti:Sapph, the light is split by a beam splitter such that a part of the light is directly used as the pump input and a part is guided to an optical parametric oscillator (OPO) to generate the input pulses at a central wavelength of $\lambda_{in}=1553.39$ nm. Following the OPO, the input pulses are sent through a neutral density filter (ND) to reduce the power to the single-photon level. After this, the light is guided to a commercial waveshaper to be able to shape the input pulses. Then, they are coupled out of the fiber to free space and pass through a half waveplate to the set polarization.

To shape the temporal mode of the pump, the pump pulses from the Ti:Sapph are sent through a home-built pulseshaper in the form of a 4-f-line with a spatial light modulator (SLM) for shaping the pump spectrum [189]. Afterwards, the pump pulses are guided to a half waveplate by a pick up mirror. The input and pump pulses are combined by a dichroic mirror (DM) and then coupled by a single aspheric lens to the periodically poled TFLN sample which is placed on a home-built oven for temperature stabilization. After the sample, the transmitted light is collected by a lens. The generated SFG light is separated from the input and pump light by a dichroic mirror and coupled to a fiber which is either connected to an avalanche photodiode (APD) or a single-photon sensitive spectrometer consisting of a spectrograph combined with a CCD camera. The input light is separated by a second dichroic mirror and also fiber coupled and connected to an SNSPD.

6.4 Benchmarking of a Quantum Pulse Gate

To benchmark our TFLN QPG, we first have to check, if the process is group velocity matched which results in a horizontal phase matching function. Therefore, we map the complete phase matching function of our device. Afterwards, we evaluate how well the device can distinguish between different temporal modes by measuring the temporal mode selectivity. Lastly, the conversion efficiency of the device is determined.

6.4.1 Characterization of the Dispersion

To check if the group velocity matching condition of our QPG is achieved, we map the phase matching function. For this, we use the same method as in section 5.2.4. This means, we slightly modify the setup shown in Figure 6.2 by bypassing the home-built waveshaper of the pump pulses and using a tuneable CW-laser in the telecom regime as input light.

We set the pump wavelength to 865 nm and scan the CW-laser from 1520 nm to 1580 nm in 0.2 nm steps. For each wavelength step, we measure the generated SFG spectrum with a single-photon sensitive spectrometer. Then, we change the pump wavelength by 5 nm and repeat the scan with the CW-laser. We scan the pump wavelength from 865 nm to 885 nm in 5 nm steps. In the end, we combine all measured SFG spectra to retrieve the phase matching function. We perform this measurement for multiple waveguides and poling periods and show here the waveguide with a poling period of $3.178 \mu\text{m}$ since this waveguide exhibits a phase matching close to our expectation.

The measured phase matching function is close to horizontal which indicates that we nearly reached group velocity matching between the input and pump wavelength (see Figure 6.3). The phase matching function shows a brighter spot at 1550 nm and 559 nm due to the

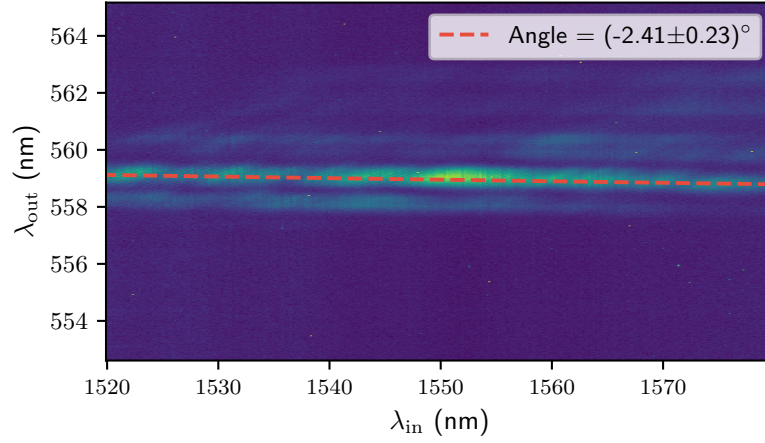


Fig. 6.3: Measured phase matching function of the QPG in TFLN. The close to horizontal phase matching function indicates that we are close to group velocity matching between the pump and input wavelength.

wavelength dependent alignment. Moreover, we can observe the expected sinc²-like shape of the phase matching function (see Equation 2.115). We can observe more pronounced side peaks at higher wavelengths which can be due to inhomogeneities in the waveguide geometry [148].

The FWHM of the measured phase matching function is 0.48 nm and the measured phase matching angle is $(-2.41 \pm 0.23)^\circ$. Both slightly deviate from the simulation which might be due to a minor different fabricated waveguide geometry. But even with the slightly negative phase matching angle, we can still use the device as a QPG. Thus, in the next step, we will check the performance characteristic of our device.

6.4.2 Temporal Mode Selectivity

To show the mode selective functionality of our QPG, we evaluate the temporal mode selectivity as defined in Equation 2.144. The temporal mode selectivity is, as discussed in subsection 2.5.3, a figure of merit which determines how well the device can distinguish between a specific temporal mode from other orthogonal temporal modes. To measure the temporal mode selectivity, we use the setup shown in Figure 6.2. We use both pulseshapers to shape the pump and input pulses in HG modes. The fundamental HG₀ mode has a bandwidth of 1 THz FWHM in amplitude. We can test the temporal mode selective behavior with these measurements since we probe it with input and pump mode with the same shape and with different shapes.

To determine the temporal mode selectivity, we sequentially shape the input and pump pulse in HG modes up to the fourth order. For each mode combination, we measure the spectrum of the generated output light with the single-photon sensitive spectrometer. For every measured spectrum, we first subtract a background noise consisting of dark counts and electronic noise from the spectrometer. For this, we record a spectrum of the spectrometer with the same settings as we use for the measurements but with all lasers turned off. Afterwards, we apply a 0.1 THz frequency filter centered at the main peak of the output spectrum to eliminate the side lobes. The filter bandwidth is chosen for easy comparison with other works. As a final step, we integrate over the remaining area to retrieve the counts to

6.4. Benchmarking of a Quantum Pulse Gate

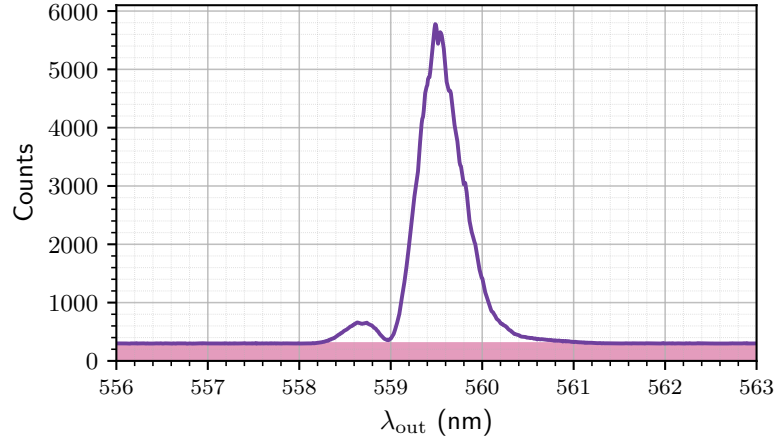


Fig. 6.4: Measured counts from the single-photon sensitive spectrometer as a function of the output wavelength. We show the measured spectrum from which we subtract the pink shaded noise floor. Moreover, the vertical shade area displays the applied frequency filter.

calculate the temporal mode selectivity (see Figure 6.4). To do so, we evaluate for each input mode the ratio of correct counts (counts from the measurement with same mode order of the input and pump mode) to the total number of counts (sum of counts of all mode combinations).

We measure the temporal mode selectivity first for high-photon numbers since the temporal mode selectivity should be independent of the number of photons due to the fact that the modal properties do not depend on the intensity. Afterwards, we also measure it at low-photon numbers, since the device should be used at the single-photon level in the end. For the measurement at high-photon numbers, we use a pump power of 2 mW and an input

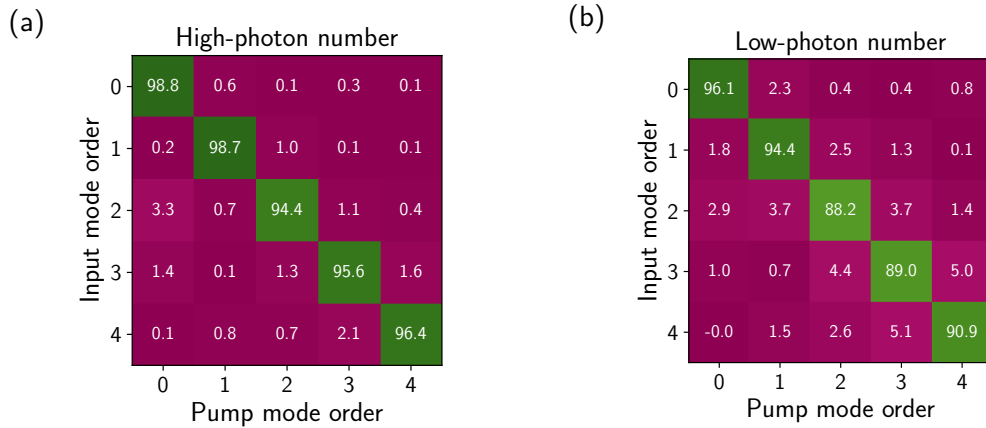


Fig. 6.5: Temporal mode selectivity plots for the first five HG modes. (a) For a high photon number we reach an average temporal mode selectivity of $(96.8 \pm 1.7) \%$. (b) For a low photon number the average temporal mode selectivity is $(91.7 \pm 3.1) \%$.

Chapter 6. Quantum Pulse Gate in TFLN

pulse with $3 \cdot 10^4$ photons per pulse, both measured before the waveguide. We show the temporal mode selectivity as a function of the input and pump mode order in Figure 6.5 (a). We can observe that we measure a high temporal mode selectivity of on average of $(96.8 \pm 1.7) \%$ which shows that our TFLN QPG can very well distinguish between different temporal modes.

To test the device for quantum applications, we reduce the photon number of the input pulse to 1 photon per pulse on average. The result is shown in Figure 6.5 (b). Here, we can again see that we measure a high model contrast of $(91.7 \pm 3.1) \%$ on average which is slightly reduced compared to the measurement with a high-photon number. The reduction might be due to noise for example from the detector or from the nonlinear process, but this needs to be investigated in more detail. But, overall, also at the single-photon level the device performance excellently and similar to standard devices [81, 178, 186].

To conclude, for both measurement settings, we can see a distinct mode selective behavior. Moreover, the measured temporal mode selectivity is comparable with that of QPGs realized in titanium indiffused lithium niobate waveguides showing that we realized a device of high standard.

Lastly, we want to investigate if our QPG in TFLN exhibits an enhanced conversion efficiency.

6.4.3 Conversion Efficiency

As the last step, we measure the conversion efficiency of our device to check if our TFLN QPG is more efficient than a QPG in titanium indiffused lithium niobate waveguides. For this, we use the same setup as before. This time, we shape both pulses, the input and pump pulse, as HG₀ modes with a bandwidth of 1 THz. To estimate the conversion efficiency, we measure the depletion of the input light. We use this method instead of measuring the generated output light since by measuring the depletion we can directly estimate the conversion efficiency. To evaluate the conversion efficiency from the measured generated output light requires the knowledge about all losses of the output light. To estimate the conversion efficiency, we measure the transmitted input light with a SNSPD after the waveguide in two settings. First, we measure the counts of the transmitted input light when the input and pump laser are turned on (C_{open}) such that conversion takes place. Second, we turn off the pump laser which means no conversion is happening anymore and we measure the complete transmitted input light passing through the waveguide. The ratio between the two measurements reveals the portion of the input light which is converted. Thus, we can evaluate the internal conversion efficiency via the following equation [185]:

$$\eta = 1 - \frac{C_{\text{open}}}{C_{\text{blocked}}}. \quad (6.1)$$

We perform this measurement for different pump powers and measure the pump power before the waveguide. We use the pump power before the waveguide since it is challenging to estimate the coupling efficiency when coupling two distinct wavelengths with one single lens into a TFLN waveguide. This means the pump power is considerable lower inside the waveguide which means we underestimate drastically the internal conversion efficiency. We plot the conversion efficiency as a function of the pump power before the waveguide in Figure 6.6. Here, we subtract background and pump-induced noise by taking a measurement with the input laser blocked. The measured conversion efficiency increases from $(15.5 \pm 0.2) \%$ at a pump power of 1 mW up to $(89.6 \pm 0.1) \%$ at a pump power of 45 mW. This value is the highest experimentally measured value beating the highest literature value

6.4. Benchmarking of a Quantum Pulse Gate

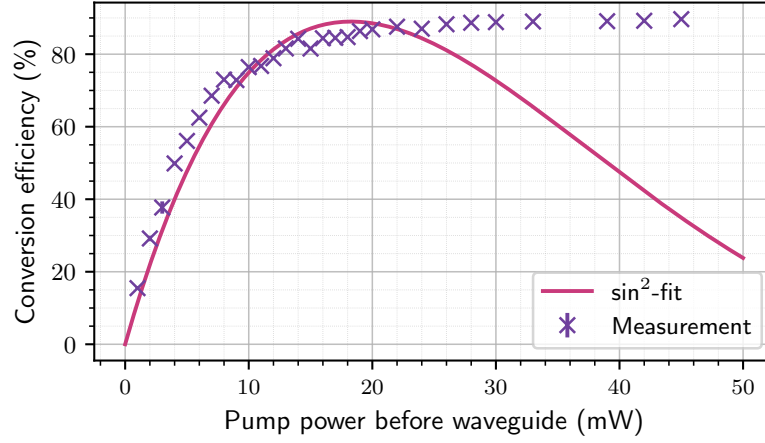


Fig. 6.6: Conversion efficiency as a function of the pump power before the waveguide. We also plot a $\sin^2$ -fit to show the expected behavior of the conversion efficiency. A possible explanation for the deviation between the fit and the measurement might be time-ordering effects.

of 87.7 % at a lower pump power due to the increased poling length [81]. The conversion efficiency stays constant after a pump power of around 15 mW. Usually, we would expect that the conversion efficiency has a $\sin^2$ -like trend which we also plot in Figure 6.6 (magenta line). This means that the efficiency increases until all the incoming photons are converted. After this point, the reverse process happens and the conversion efficiency decreases again. In our case the conversion efficiency stays constant which we can explain with time-ordering effects [88] which needs to be investigated in more detail.

To estimate the normalized conversion efficiency, we apply a fit according to Equation 2.145. The high conversion efficiency at moderate pump powers leads to a normalized conversion efficiency of $1115 \text{ W}^{-1} \text{ cm}^{-2}$. The normalized conversion efficiency is considerable underestimated since we assume that the complete pump power before the waveguide is coupled to the waveguide. Thus, this value provides only a lower bound. However, even the lower bound is already three order of magnitude higher than the best value reported for a QPG in titanium indiffused waveguides in lithium niobate with the same wavelength combination [89]. Approaches based on type 0 conversion in an approximate group-velocity matching configuration at different wavelengths have shown high conversion efficiencies [190, 191], but more detailed investigations are required for examining their noise performance at the single-photon level. A rough estimation of the coupling efficiency by simulating the overlap of the waveguide mode to the mode distribution given by the lens yields a coupling efficiency of around 28 %. Assuming that we couple 28 % of the pump power before the waveguide into the waveguide leads to a normalized conversion efficiency of $6000 \text{ W}^{-1} \text{ cm}^{-2}$.

This enhancement of the normalized conversion efficiency has two reasons. First, TFLN waveguides have an increased mode overlap between different frequency modes due to the high-index contrast waveguides compared to titanium indiffused waveguides in lithium niobate. The mode overlap, see Equation 2.108, of the employed structure is in the order of $1.7 \mu\text{m}^{-2}$. For a QPG in titanium indiffused waveguide the mode overlap is the order of $0.01 \mu\text{m}^{-2}$ which results in an increase of two orders of magnitude [188]. Second, TFLN allows for dispersion engineering which enables us to utilize a type 0 SFG process with the d-tensor element d_{33} instead of a type 2 process which is based on the d-tensor element

Chapter 6. Quantum Pulse Gate in TFLN

d_{31} . This leads to an increase of the efficiency of one order of magnitude. Thus, due to the higher d-tensor element of the nonlinear process and the increased mode overlap, the normalized conversion efficiency is enhanced by three order of magnitude.

This result highlights our outstanding device performance and opens up the possibility to use QPGs in large scale quantum experiments, such as quantum key distribution or quantum computing schemes, due to the enhanced normalized conversion efficiency. Moreover, these results underlines our capabilities of designing, fabricating and characterizing dispersion engineered quantum devices in TFLN.

Summary

In this chapter, we discussed the design and characterization of the first QPG in TFLN. We started with dispersion engineering to find a suitable waveguide geometry which provides group velocity matching. Afterwards, we explained the setup used to benchmark our device. Finally, we characterize the QPG by first determining the dispersion by mapping the phase matching function. Then, we check the selective mode behavior of the device by evaluating the temporal mode selectivity which showed a distinct mode selective behavior with an average temporal mode selectivity of $(96.8 \pm 1.7) \%$ for high-photon numbers and of $(91.7 \pm 3.1) \%$ for low-photon numbers inline with standard QPGs. Lastly, we measure a conversion efficiency of maximally $(89.6 \pm 0.1) \%$ which is the highest reported value. This leads to a lower bound of the normalized conversion efficiency of $1115 \text{ W}^{-1} \text{ cm}^{-2}$ which shows that our TFLN device outperforms a standard QPG by three orders of magnitude. This result paves the way for the implementation of QPGs in large scale quantum networks. With this work, we successfully showed our approach starting from device design to fabrication and lastly to quantum characterization to demonstrate dispersion engineered quantum devices in TFLN.

7

Conclusion and Outlook

7.1 Conclusion

The aim of this thesis was to harness the advantages of TFLN for improved dispersion-engineered quantum devices. For this, we started with the basics namely the characterization of low-loss waveguides, moved on to passive quantum effects and finally demonstrated dispersion engineered active quantum devices. This approach enables us to reliably design, fabricate and characterize quantum devices in TFLN.

We started with the basic building block of integrated optics which are low-loss waveguides. These are especially crucial for integrated quantum optics due to the no cloning theorem. We compared two common methods to estimate the propagation losses of waveguides namely the Fabry-Pérot method and the racetrack resonator method. By measuring and comparing three samples with different top widths ranging from single-mode to multi-mode waveguides, we concluded that the Fabry-Pérot method can only be utilized when investigating the end face reflectivity in depth. With the racetrack resonator method, we showed that the propagation losses decay with increasing top width from (1.462 ± 0.195) dB/cm for a top width of $1 \mu\text{m}$ to (0.024 ± 0.036) dB/cm for a $3 \mu\text{m}$ top width. This is due to the fact that the mode is less confined in the waveguide for broader top widths leading to less overlap with the rough sidewalls. Thus, when designing an optical circuit, broader waveguides should be used to route the light while smaller waveguides are beneficial at interaction regions such as poled areas where a single mode operation is crucial to avoid higher order processes. Moreover, we showed that our best losses of (2.4 ± 3.6) dB/m for multi-mode waveguides are in the same order of magnitude compared to the state of the art multi-mode waveguides. This indicates that we are able to fabricate and characterize low-loss waveguides in TFLN which are highly relevant for quantum applications.

To demonstrate that TFLN is a highly interesting material platform for integrated quantum optics, we moved on to a genuine quantum effect. We showed the most fundamental quantum phenomenon which is the Hong-Ou-Mandel interference in a photonic integrated circuit in TFLN. This interference between two single photons in a balanced beam splitter is crucial and opens up the possibility towards the integration of complex quantum circuits. To show this interference on chip, we simulated, designed and characterized directional couplers in TFLN. We showed that the simulation of the directional coupler matched with our experiments. Moreover, we measured a HOM dip with a visibility of $(93.5 \pm 0.7) \%$ which is reduced due to the source visibility and the non-perfect 50:50 splitting of the

Chapter 7. Conclusion and Outlook

directional coupler. However, we still have a 3 % difference between the measured and maximal achievable visibility which might be due hybrid modes. The results show that we are able to design, fabricate and characterize directional couplers in TFLN which are a main building block for various integrated devices.

To use TFLN as material for integrated quantum optics, we want to be able to utilize nonlinear processes and exploit the possibility of dispersion engineering in TFLN waveguides. Thus, we focused on the generation of single photons with tailored properties via parametric down-conversion in dispersion engineered TFLN waveguides. First, we investigated the phase velocities to realize a PDC source generating photons at two distinct wavelengths. Based on our design, we realized a highly non-degenerate PDC source in TFLN generating photons at 815 nm and 1550 nm. We experimentally characterized the spectral and quantum properties of our source in depth. With the results, we were able to verify our simulations showing that they are highly suitable for the design of such sources. We measured a heralded second-order correlation function $g_h^{(2)}(0)$ of $(6.7 \pm 1.1) \cdot 10^{-3}$ showing that our source generates single photons with a high quality. Moreover, we determined a brightness of the source of $0.44 \cdot 10^7 \frac{\text{pairs}}{\text{s} \cdot \text{mW} \cdot \text{GHz}}$ which is on par with other sources in TFLN. Furthermore, we started to discuss the topic of waveguide fabrication tolerances. We explored this topic in more detail and investigated the influence of thin-film thickness variations. A varying thin-film thickness distorts the phase matching function. We simulated how an adaptive top width can be used to counteract the influence of the thin-film variations and reconstruct the phase matching function. Second, we dived deeper into the topic of dispersion engineering and discussed how PDC sources with symmetric group velocity matching can be realized which are highly interesting for various quantum optical experiments. Here, we examined two different approaches. The first one is to increase the temperature of the device to achieve the necessary dispersion and the second approach is to utilize suspended waveguides. This chapter showed on the one hand, the great consistency between our simulation and experiment and on the other hand, the capability of dispersion engineering in TFLN waveguides.

To show advanced and improved dispersion engineered devices in TFLN, we realized the first TFLN quantum pulse gate. First, we designed the waveguide geometry such that we achieve group velocity matching between the input and pump wavelength. Then, we experimentally characterized the QPG by evaluating the temporal mode selectivity. Here, we showed a clear mode selective behavior with an average temporal mode selectivity of $(96.8 \pm 1.7) \%$ for high-photon numbers and of $(91.7 \pm 3.1) \%$ for low-photon numbers. The measured visibility is on par with the temporal mode selectivity of QPG realized in standard lithium niobate and thus, showing our great device performance. Lastly, we evaluate the conversion efficiency of the QPG in TFLN. We measure a conversion efficiency of $(89.6 \pm 0.1) \%$ which is higher than the highest reported value in literature so far. This also leads to a normalized conversion efficiency of $1115 \text{ W}^{-1} \text{ cm}^{-2}$ which is three orders of magnitude higher compared to QPGs in titanium indiffused waveguides in lithium niobate. This high conversion efficiency leads to time ordering effects which need to be investigated in depth in the future. This high conversion efficiency sets the stage for the implantation of QPGs in large scale quantum networks. This demonstrates that TFLN can outperform devices based on titanium indiffused waveguides and that we can realize highly efficient dispersion engineered devices in TFLN.

In conclusion, during the scope of this thesis we gained profound knowledge about the design

and characterization of linear and quantum devices in TFLN. This lays a strong foundation for further research of integrated quantum optics in TFLN waveguides in our group.

7.2 Outlook

The demonstrated results of this thesis pave the way for further continuation of dispersion engineered quantum devices based on TFLN.

A first valuable step forward could be to experimentally verify that the top width can be used to counteract thin-film thickness variations. For this, the thin-film thickness is measured with an ellipsometer and we design a sample with one waveguide with a constant top width and the second waveguide with an adapted top width matched to the thin-film thickness. With this sample, we can simulate the phase matching of the varying thin-film and compare it to the measurements to validate the simulations. This technique allows us to realize nonlinear devices with longer poled areas leading to a higher efficiency and smaller bandwidth.

Another important step for the community would be to experimentally demonstrate the proposed waveguide designs for the symmetric group velocity matched PDC sources. With this, we would not only demonstrate an important source for many quantum optical applications, but we can also use the measurements to validate our simulations of the TM modes. During the scope of this thesis, we investigated experimentally only type 0 processes for which all involving fields are TE polarized. Thus, by measuring a type 2 PDC source, we can check the simulations for TM modes.

On a broader scope, we should use the gained knowledge about PDC sources and directional couplers and combine both on chip to achieve pump suppression. After the generation of the signal and idler photons, pump light is still propagation with the generated photons in the waveguide which for example blinds the SNSPDs when taking measurements at the single photon level. To get ride of the pump light, we can use a directional coupler to separate the pump light from the signal and idler photons. The pump light would stay in the initial waveguide and the signal and idler photons would couple to the second waveguide of the directional coupler.

We can expand our horizon even more by combining a PDC source with a highly efficient QPG in TFLN. The generated modes of a highly multi-mode type 0 PDC enter a QPG which then selectively performance operations on the modes in a full programmable way [176] meaning we can realize frequency based Gaussian boson sampling. The necessary pulse shaper for the QPG operation can also be implemented in TFLN [192, 193] such that we can achieve a high-dimensional temporal-mode based quantum process on a single chip.

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A

Appendix

A.1 Patents

Phasenangepasste Kompensation in Wellenleitern

Silia Babel, Laura Padberg and Christine Silberhorn, filed on 31/07/2024, Patent number DE 10 2024 121 775.8

Method for manufacturing a photon emitter

Silia Babel, Laura Padberg and Christine Silberhorn, filed on 30/07/2025, Patent number 19/285,966

Photonenpaarquelle

Silia Babel, Laura Padberg and Christine Silberhorn, filed on 27/11/2025 Patent number DE 10 2025 149 426.6

A.2 Publications

First author publications

Demonstration of Hong-Ou-Mandel interference in an LNOI directional coupler

Silia Babel, Laura Bollmers, Marcello Massaro, Kai Hong Luo, Michael Stefszky, Federico Pegoraro, Philip Held, Harald Herrmann, Benjamin Brecht, Christof Eigner, Laura Padberg, and Christine Silberhorn
Optics Express, 31(14), 23140-23148. (2023)

We have demonstrated a Hong-Ou-Mandel interference between two single photons in a directional coupler in TFLN. I designed and characterized the directional coupler, realized a new setup and measured the Hong-Ou-Mandel dip.

Ultrabright, two-color photon pair source based on thin-film lithium niobate for bridging visible and telecom wavelengths

Silia Babel, Laura Bollmers, Franz Roeder, Werner Ridder, Christian Golla, Ronja Köthemann, Bernhard Reineke, Harald Herrmann, Benjamin Brecht, Christof Eigner, Laura Padberg, and

Appendix A. Appendix

Christine Silberhorn

Optics Express, 33(25), 52729-52743. (2025)

This work presents the first realization of a highly non-degenerate PDC source in TFLN. Here, we generate one photon around 1550 nm and the other photon at around 815 nm making this source highly interesting for many applications such as the use of it as a heralded single photon source or for quantum spectroscopy. I designed the PDC source and performed the in-depth characterization.

Demonstration of a quantum pulse gate in thin-film lithium niobate

Silia Babel*, Alejandra Alarcón*, Laura Serino, Christian Golla, Laura Bollmers, Sebastian Lengeling, Jiayu Yang, Bernhard Reineke, Christof Eigner, Benjamin Brecht, Laura Padberg, Marko Loncar, and Christine Silberhorn

* These authors contributed equally to this work.

In preparation

We have shown the first quantum pulse gate in TFLN. I designed the samples as well as built a new setup for the measurements. Moreover, I did the initial characterization of the sample and took the first measurements showing the performance of the device.

Next to these main projects, I also contributed to the following works.

Author contributions

Estimation of losses caused by sidewall roughness in thin-film lithium niobate rib and strip waveguides

Manfred Hammer, Silia Babel, Henna Farheen, Laura Padberg, J. Christoph Scheytt, Christine Silberhorn, and Jens Förstner

Optics Express, 32(13), 22878-22891. (2024)

I characterized the waveguides and provided the measurement results.

Spectro-temporally tailored Non-Gaussian Quantum Operations in Thin-Film Waveguides

Peter Namdar, Patrick Folge, Carlos E. Lopetegui, Silia Babel, Benjamin Brecht, Christine Silberhorn, and Valentina Parigi

arXiv preprint arXiv:2508.04578. (2025)

I provided information about waveguide designs in TFLN waveguides and discussed simulation results.

Fabrication tolerances in TFLN waveguides

Jonas Babai-Hemati, Ernst-Lukas Kuhlmann, Silia Babel, Noah Spiegelberg, Christian Golla, Bernhard Reineke, Michael Rüsing, Laura Bollmers, Christof Eigner, Laura Padberg, Harald Herrmann, and Christine Silberhorn

In preparation

A.3. Conference Contribution

I helped with the initial idea of the paper as well as provided simulation data. Moreover, I simulated first results.

High-Performance Photon-Pair Source Using a Two-Photon-Polymerized Strip-Loaded Lithium Niobate Waveguide

Muhamed A. Sewidan, Alexandra Rittmeier, Laura Bollmers, Silia Babel, Elisavet Chatzizyrlı, Juan S.S. Duran Gomez, Laura Padberg, Christof Eigner, Christine Silberhorn, Douglas Bremner, Andreas Wienke, Dietmar Kracht, Moritz Hinkelmann, and Michael Kues
In review at Optics Express

I provided knowledge about simulation of TFLN waveguides.

Surface Acoustic Wave Mediated Brillouin Temperature Sensor on Thin Film Lithium Niobate

Varun Marathan Kovil, Justus Nölkenmeier, Pratham Kulkarni, Gladys Jara-Schulz, Silia Babel, Christian Golla, Rajkumar Jadhav, Christof Eigner, Laura Padberg, Christine Silberhorn, and Birgit Stiller
In preparation

I designed waveguide masks for fabrication and characterized the waveguides in terms of losses.

A.3 Conference Contribution

- DPG spring meeting, Erlangen (online), 2022, talk
- International Workshop on Polar Oxides: Lithium Niobate and Related Compounds, Goslar, 2022, talk
- SFB/TRR 142 meeting, Bad Sassendorf, 2022, talk
- DPG spring meeting, Hannover, 2023, talk
- CLEO US, San Jose, 2023, talk
- PhoQS Symposium Quantum Spots, Bad Driburg, 2023, poster
- CLEO US, Charlotte, 2024, talk
- Quantum Photonics Spotlight, Paderborn, 2024, poster
- Research Visit Harvard University, Boston, 2024, talk