Quantum sensor using guided matter waves in optical fibers

Xin, Mingjie

2018

Xin, H. (2018). Quantum sensor using guided matter waves in optical fibers. Doctoral thesis, Nanyang Technological University, Singapore.

<https://hdl.handle.net/10356/89641>

<https://doi.org/10.32657/10220/47709>

**NANYANG
TECHNOLOGICAL
UNIVERSITY**

SINGAPORE

QUANTUM SENSOR USING GUIDED MATTER WAVES IN OPTICAL FIBERS

XIN MINGJIE

SCHOOL OF PHYSICAL AND MATHEMATICAL SCIENCES

2018

QUANTUM SENSOR USING GUIDED MATTER WAVES IN OPTICAL FIBERS

XIN MINGJIE

School of Physical and Mathematical Sciences

A thesis submitted to the Nanyang Technological University
in partial fulfillment of the requirements for the degree of
Doctor of Philosophy

2018

© Copyright by Xin Mingjie 2018

All Rights Reserved

Statement of Originality

I hereby certify that the work embodied in this thesis is the result of original research done by me except where otherwise stated in this thesis. The thesis work has not been submitted for a degree or professional qualification to any other university or institution. I declare that this thesis is written by myself and is free of plagiarism and of sufficient grammatical clarity to be examined. I confirm that the investigations were conducted in accord with the ethics policies and integrity standards of Nanyang Technological University and that the research data are presented honestly and without prejudice.

19/02/2019

.....
Date

.....
Xin Mingjie

Supervisor Declaration Statement

I have reviewed the content and presentation style of this thesis and declare it of sufficient grammatical clarity to be examined. To the best of my knowledge, the thesis is free of plagiarism and the research and writing are those of the candidate's except as acknowledged in the Author Attribution Statement. I confirm that the investigations were conducted in accord with the ethics policies and integrity standards of Nanyang Technological University and that the research data are presented honestly and without prejudice.

19/02/2019

.....
Date

.....
Lan Shau-Yu

Authorship Attribution Statement

This thesis contains material from one paper published in the following peer-reviewed journal where I was the first author.

Chapter 6 is published as M. Xin, W. S. Leong, Z. Chen, S.-Y. Lan. An atom interferometer inside a hollow-core photonic crystal fiber. *Sci. Adv.* 4, e1701723 (2018). DOI: 10.1126/sciadv.1701723.

The contributions of the co-authors are as follows:

- A/Prof Lan Shau-Yu provided the initial project direction and contributed to designing the experiment, analyzing the data, and preparing the manuscript.
- Mr Leong Wui Seng assisted in the collection of the data and contributed to designing the experiment, analyzing the data, and preparing the manuscript.
- Dr Chen Zilong I contributed to designing the experiment, analyzing the data, and preparing the manuscript.

19/02/2019

.....
Date

.....
Xin Mingjie

Abstract

Coherent interaction between matter waves and light is the core of quantum sciences and technologies. In the optical domain, coherent manipulation of quantum states of atoms using light has been employed in applications such as precision metrology, quantum simulation, quantum computing, etc. Miniaturizing technologies of atom-light-based quantum systems for practical applications has been a longstanding goal since their development. However, the diffraction nature of light has limited the scalability of the quantum systems. Highly sensitive light-pulse atom interferometers in free space are tens of centimeters in size and, therefore, they are hardly being used in micro-scale systems for wide applications.

In this work, we optically trap cold ^{85}Rb atoms inside a hollow-core photonic crystal fiber and use the optical waveguide mode inside the micro-scale structure to manipulate atoms as beamsplitter and mirror pulses to form an inertia-sensitive atom interferometer which can be used to measure the gravitational acceleration g . Further, the atom interferometer is not limited by the diffraction nature of the light. Coherent interaction between a waveguide mode and a quantum system, could lead to constructing mobile and miniature quantum hybrid platforms for precision measurement and quantum sensing. This work is related to and can be applied to a wide range of disciplines, including photonics, fundamental physics, and quantum sensor technology.

Moreover, the dephasing mechanism of quantum superposition states of the atoms and the transportation of quantum superposition inside the fiber are studied. We demonstrate coherent guiding of ground-state superpositions of ^{85}Rb atoms over a centimeter range and over hundreds of milliseconds inside a hollow-core fibre. The transportation distance is limited by the inhomogeneity of ambient magnetic field in our apparatus and not by the fiber structure. Our experiment establishes an important step towards a versatile platform that can bring long-lived quantum spin coherence outside the traditional bulky apparatus and over variable distances. This can lead to applications in quantum information networks and matter wave circuit for quantum sensing.

Acknowledgement

My Ph.D. work wouldn't be possible without the encouragement, support, help and assistance from lots of nice people I overlapped during my whole graduate career.

My thanks go first and foremost to my supervisor Prof. Lan Shau-Yu for his guidance, insight, support, and encouragement during my whole Ph.D. study. It is a great honor to work closely with him. His scientific attitude, deep knowledge, passion, and patience make a profound impression to me. I benefited a lot from him on the research and obtained an efficient and professional training in the atomic physics field. He also provided lots of opportunities for me to attend the international conferences on the atomic physics. I feel grateful and lucky for the opportunity to join his research group and work on this interesting project.

I wish to thank my thesis committee members, Prof. Kim Young-Jin and Prof. Rainer Helmut Dumke for their comments, discussion, suggestions and support on the progress of my research. I also would like to thank them for taking the time to read my research report and give invaluable comments. I also would like to thank Prof. Xiong Qihua and Prof. Sun Handong for their help and advice in the application of my Ph.D. programme and their support in my qualification examination.

I would like to thank all members in our group, Dr. Roy Arpan, Dr. Kuan Pei-Chen, Dr. Chen Zilong, Mr. Chan Wei Sheng, Mr. Leong Wui Seng, Mr. Huang Chang, Mr. Sandoko Kosen, and Mr. Lim Hong Ming. I am very lucky to work with them in this group.

Dr. Roy Arpan and Mr. Chan Wei Sheng are kind enough to take me under their experimental wings in the beginning of my Ph.D. study, I learned lots of knowledge and experimental skills from each one of them. I am grateful to work closely with these two smart guys. I would like to thank Dr. Chen Zilong, who is a good leader for our project. He is very professional and has lots of experience on the atomic and optical physics. His guidance, advise and insight is very helpful for this research. For the stimulating discussion with

Dr. Kuan Pei-Chen, I learned lots of knowledge and experimental techniques. I also enriched my knowledge in other research areas by the daily interactions with him and he made a lot of fun in the “NTU GGland”.

In particular, I would like to thank Mr. Leong Wui Seng, who worked with me on this hollow-core fiber project. Wui Seng is an extraordinary partner in research and a nice friend in my daily life. This project would hardly be possible without his generous assistance and patient help. He also gave me many useful and helpful suggestions in writing this thesis.

It is a great honor to have collaboration with Mr. Lim Hong Ming, Mr. Huang Chang, Mr. Sandoko Kosen, and Mr. Ian Yap Chang Jie during various stages of my four-year-long graduate career. They provided lots of help on my experiment and I greatly enjoyed the collaboration with each one of them.

Furthermore, I want to thank all my friends for their encouragement. Especially, I wish to thank Xu Yang for being a special person in my life. Thanks to her for the love, patience, and understanding.

Finally, I would like to thank my family, my grandparents and my parents Xin Fengkun and Wu Liqing, who always support me throughout my life. Their unconditional love and support made me get through the hard four-year-long graduate career.

*Dedicated to my parents Xin Fengkun and Wu Liqing,
and my grandparents.*

Contents

Abstract	i
Acknowledgement	ii
Table of Contents	vii
1 Introduction	1
1.1 Motivation	1
1.2 An atom interferometer inside a hollow-core photonic crystal fiber	4
1.3 Transporting long-lived quantum spin coherence inside a hollow-core fiber	6
1.4 Overview	8
2 Laser cooling and trapping	9
2.1 Atomic structure	9
2.2 Atom-light interaction	11
2.3 Optical forces	14
2.3.1 Scattering force	14
2.3.2 Dipole force	16
2.4 Laser cooling	18
2.4.1 Doppler cooling and optical molasses	18
2.4.2 Sisyphus cooling	21
2.5 Magneto-optical trap (MOT)	23
2.5.1 Trap loading	23
2.5.2 Number of atoms	25
2.5.3 Temperature measurement	28
3 Atom interferometry	31
3.1 Ramsey interferometer	31

CONTENTS

3.2	Light-pulse atom interferometry	33
3.2.1	Raman Mach-Zehnder interferometer	33
3.2.2	Interferometer phase shifts	36
3.3	The path integral approach	37
3.3.1	Propagation phase	38
3.3.2	Separation phase	40
3.3.3	Laser interaction phase	40
3.4	Precision Gravity Measurements	42
4	Experimental apparatus	45
4.1	Fiber selection	45
4.2	Vacuum system	47
4.3	Laser system	50
4.3.1	Cooling lasers	51
4.3.2	State preparation and detection Lasers	52
4.3.3	Dipole laser	53
4.3.4	Raman Lasers	55
4.3.5	Optical pumping laser	56
4.4	Laser cooled atomic sources	57
4.5	Control system	59
4.6	Microwave frequency generation	61
4.7	Detection devices	62
5	Cold atoms inside a hollow-core optical fiber	65
5.1	Fiber preparation	65
5.2	Atoms inside the fiber	67
5.3	Number of atoms inside fiber	72
6	Atom interferometry inside a hollow-core photonic crystal fiber	75
6.1	Temperature of the atoms inside the fiber	76
6.2	Atom interferometer inside fiber	79
6.3	Temperature of atoms in the atom interferometer	86
7	Transporting long-lived quantum spin coherence inside fiber	91
7.1	Atom loading using optical lattice	93
7.2	Lifetime of atoms inside the fiber	97
7.3	Temperature measurements of atoms inside the fiber	98
7.4	Coherence time of atoms inside fiber	99

CONTENTS

7.5 Transportation of quantum superposition	107
8 Conclusions	111
8.1 Summary	111
8.2 Outlooks	112
Appendix	115
Bibliography	123

Chapter 1

Introduction

1.1 Motivation

Precision measurements with atomic systems are used to study fundamental physics as well as applied sciences. In particular, over the past two decades, matter waves interferometers have played a significant role in testing general relativity [1], measurements of the fine-structure constant [2,3], and Newtonian gravitational constant G [4,5]. They have also been used as precise inertial sensors to measure gravitational acceleration [6–8], gravity gradient [9], gravity curvature [10] and rotation [11]. In light-pulse atom interferometers, atoms are spatially split, reflect, and recombine through two separated paths by atom-light interactions to form interference fringes. The phase difference from the evolution of the wave-function along the two distinct trajectories can be determined through interference signals, which is sensitive to the potential field applied to the atoms [12,13].

Although free space atom interferometers have reached unprecedented sensitivity, thanks in part to the development of large momentum transfer beam splitters [14] and studies of noise sources and systematic effects [15], all high sensitive atom interferometers utilize atom-photon interactions to split, recombine, and interfere matter waves in free space [1–11]. This provides a straightforward experimental design and has the capability to work with large atom numbers. However, due to the diffraction nature of light, the diameters of the

interferometer tube are typically tens of centimeters in order to provide homogeneous light intensity throughout atoms' trajectories and avoid scattering of light on the wall of the vacuum chamber. Therefore, it is difficult to carry atoms close to source masses, and hardly to integrate them into small portable systems for wide applications. Moreover, they can't be configured to desired geometry easily. Various atomic guiding systems have been proposed and matter waves interferences have been investigated in systems, such as magnetic guide on chips and optical guide in free space. They have yet not achieved the sensitivity of atom interferometers in free space [16, 17] and they are not suitable for long distance and flexible guidance.

Alternatively, hollow-core optical fibers which could guide light over long distance can be used to guide atoms as well. Atoms can be confined by the optical dipole potential created by the optical modes in the fiber, see Fig. 1.1(a). Compared to free space optically guided matter waves, the optical beams guided in the fiber is in principle diffraction free which allows homogeneous atom-light interaction strength over long distances, see Fig. 1.1(b). The sensitivity of light-pulse atom interferometers scales with the momentum of light imparted on atoms during pulses. Multi-order Bragg diffraction atom interferometers have been demonstrated with 24-photon-momentum transfer beam splitters [14]. However it was mainly limited by the optical power. The strong confinement of optical beams inside the optical fiber could lead to studies of high-order Bragg Diffraction with small optical power. The coherence of optical field in spatial and time domains are crucial for coherence manipulation of quantum states so the wave front distortion of optical beam is another limitation to the sensitivity of atom interferometers, see Fig. 1.1(c). The small aperture of fibers serves as a wave front filter that could provide flat optical wave fronts along the fiber. The Gaussian-like optical mode in capillary fibers which is excited by grazing incidence has been used to guide the atoms [18]. However, those experiments suffered rapid attenuation of the beams guided inside the capillary and significant atom losses due to speckle and inhomogeneous optical fields which cause local absence of confinement.

Recently developed hollow-core photonic crystal fibers have the special fea-

Figure 1.1: (a) Atoms are confined in the hollow-core photonic crystal fiber by using an optical dipole trap. (b) The atom-light-interaction length inside the optical fiber is not limited by the Rayleigh length like the beams in free space. Assuming the beam waist is $5\mu\text{m}$, the Rayleigh length of the beam in free space is limited to $\sim 100\mu\text{m}$. The Rayleigh length of the beam inside fiber is only limited by the length of the fiber. (c) The wave front of the beam is normally distorted after reflection by the optical surfaces, such as mirrors and wave plate. The small aperture of the fiber serves as a wave front filter for the beams. The wave front, in principle, should be homogeneous along the fiber. (d) An example of compact atomic gravimeter based on a light-pulse interferometer [24]. The atom interferometer based on the optical fiber can be used to design more compact mobile sensors.

tures of low transmission loss and less speckle [19–21]. They could also provide a homogeneous confining potential for atoms [22, 23] over long distance. Such an optical waveguide could be used to transport atoms to the desired location for measurements. It could be used in the design of next generation compact, robust, and portable atom inertial sensors (Fig. 1.1(d)).

With the small size of the fiber core, interaction of atoms with the fiber inner wall needs to be considered. At a distance $\sim 5\mu\text{m}$ between the atoms

and the fiber wall, the dominant interaction is the Casimir-Polder force. The Casimir-Polder potential has been measured to be on the order of hundreds of mHz at this distance [25], which is far lower than the typical trapping potential of about 10 MHz used in our experiments. While the hollow-core photonic fiber is good at guiding the nearly Gaussian mode, coupling the light from free space to the fiber also could result in the excitation of higher order modes. However, these higher order modes are attenuated faster than the fundamental mode.

Another advantage of the optical dipole trap interferometer inside the optical fiber is the ability to trap the atoms close enough to the source mass for the measurement of Newtonian gravitational constant G , for example. A measurement of G using atom interferometers has been demonstrated a value of $G = 6.667 \times 10^{-11} \text{ m}^3 \text{ kg}^{-1} \text{ s}^{-2}$, estimating a statistical uncertainty of $\pm 0.011 \times 10^{-11} \text{ m}^3 \text{ kg}^{-1} \text{ s}^{-2}$, and a systematic uncertainty of $\pm 0.003 \times 10^{-11} \text{ m}^3 \text{ kg}^{-1} \text{ s}^{-2}$ [5]. That experiment was limited by the knowledge of the relative position of atoms and source mass as well as the initial position and velocity of the atoms. The fiber guided atom interferometer could provide a compact setup to improve the signal and minimize the systematic effects of the current measurement. For instance, the initial position of the atoms confined inside the fiber can be known with micrometer precision, and the distance between the atoms and source mass can be reduced to hundreds of millimeter scale.

The integration of phase coherent photonic and quantum systems inside the optical fiber has a great potential to advance the capability of atom interferometers in the field of precision measurement and quantum sensing with miniature design of apparatus and high efficiency of laser power consumption.

1.2 An atom interferometer inside a hollow-core photonic crystal fiber

Due to the diffraction nature of the light, light-pulse atom interferometers are hardly brought close to source masses or surfaces to increase the interaction signals. They are hardly being used in micro-scale systems for wide applica-

tions, such as, precision measurements with high spatial resolution or designation of micrometer-scale devices. Hollow-core photonic crystal fibers could provide a platform to guide atom with only few hundreds micrometers thickness over long distances, not limited by the diffraction. It could bring atoms based quantum sensors to locations require high spatial resolution. Over the last couple of decades, the design and manufacture of hollow-core photonic crystal fiber have been developed rapidly. The guiding mode of the fiber is a nearly single Gaussian mode and provides long distance atom-light interaction limited by the fiber transmission loss only [26, 27]. The guiding mode of the beams inside the optical fiber can be used to trap, guide, transport, and manipulate atoms. Experiments based on the hollow-core photonic crystal fiber have shown that atoms could be confined, guided, or transported using the guided modes as optical trapping beams for applications in nonlinear optics and spectroscopy [28–33]. They show the possibility of the fiber-based atom interferometers.

In this work, we develop a new experimental system to trap cold ^{85}Rb atoms into a single mode hollow-core photonic crystal fiber. The guiding mode of the optical fields inside the fiber is $22\text{ }\mu\text{m}$ in radius. Cold ^{85}Rb atoms are prepared in a magneto-optical trap during the loading process. After loading into the fiber, atoms are trapped and guided using a far red-detuned optical dipole trap with a spontaneous emission rate of few per second. This will prevent collision of atoms on the inner wall of fiber and keep them away from the surface potential. The Raman laser beams for the interferometer are coupled through both ends of the fiber.

Our demonstration shows that the coherence of quantum superposition states of atoms can be interrogated and preserved in the guiding mode of the optical fields inside the hollow-core photonic crystal fiber. We also demonstrate the gravity measurement using a free-falling atom interferometer inside the fiber. It indicates that the wave front of the optical field inside the optical fiber can serve as a phase reference for atom-light interactions. The coherent interaction between the optically guiding mode and a quantum system could lead to constructing mobile and miniature platforms for precision measure-

ments and quantum sensing. This work could provide a novel way to improve the current atom-interferometer-based applications and devices. For example, in Newtonian Gravitational constant G measurement.

1.3 Transporting long-lived quantum spin coherence inside a hollow-core fiber

One of the major tasks in modern quantum sciences and technologies is to fully exploit the potentialities of quantum mechanical systems, by guiding and transporting quantum superposition states (spin qubits) of massive particles (atoms, ions and molecules) over a controllable distance while preserving their coherence [34–36]. And long-lived coherence of quantum superposition states is one of the key constituents required for quantum network and precision measurement.

As opposed to photons, quantum superposition states of massive particles could be perturbed by the guiding environment and lose their quantum coherence. In quantum memories, a key elements in quantum internet, the coherence quantum superposition states of the massive particles are used for the information storage, which are typically located in a well-designed optical trap to minimize the decoherence. A long coherence time of quantum superposition states has been demonstrated in a one dimensional optical lattice with a “magic” magnetic field to reduce the dephasing from the inhomogeneous broadening due to differential ac Stark shift introduced by the guiding optical lattice [37]. In quantum sensing, such as atom interferometers and atomic clocks, atoms are placed in free space and decoupled from the influence of dephasing sources. So far, in free-space-based quantum systems, the coherence of quantum superposition has been pushed into a macroscopic distance regime [38]. In other cases, coherent quantum superposition of atoms is localized within a one dimension or employed single atom to suppress motional decoherence [39–41]. To extend the use of the coherent quantum superposition into three dimensions over a controllable path has been a challenging goal.

For neutral atoms, the optical traps and magnetic traps are typically used to confine the atoms in the transverse direction [42]. In the propagating direction, the transportation distance of atoms in the magnetic guide is limited by the inhomogeneous magnetic field due to imperfect design and manufacture of magnetic coils or magnets. Moreover, the magnetic field is hardly to reconfigure to the desired traveling paths of atoms [43]. A controlled quantum state transportation is demonstrated by using an optical conveyor belt in free space. Although the superposition state of atoms travels over 1 mm distance and preserves the atomic coherence over few tens of milliseconds, the guiding distance is fundamentally limited by the Rayleigh range of the optical field [44].

The optical fiber could provide a platform to extend the flexibility and scalability for guiding and transporting atoms [45, 46]. For nano-fibers, atoms are trapped in a one-dimensional optical lattice by interacting with the evanescent optical fields surrounding the surface of optical nanofibers [47, 48]. Atoms are trapped about several hundred nanometers above the nanofiber surface. The coherence properties of superposition states of atoms in the nanofiber-based dipole trap have been studied experimentally in ref. [49]. Due to the inhomogeneous broadening of differential ac stark shift from the evanescent optical potential, the coherence time of nanofiber-trapped atoms is limited to few milliseconds and will eventually be limited by the surface interaction.

Hollow-core photonic crystal fibers are another platforms for trapping and guiding atoms. They have the advantage that the distance between atoms and the inner surface of fiber is typically few to tens of micrometers away. Guiding atoms inside these hollow-core photonic crystal fibers have been demonstrated over centimeter distances [50–52]. However, the transportation of quantum superposition inside the hollow-core fiber over a long distance is still unclear and remains a challenging task. A free-falling Mach-Zehnder atom interferometer has been demonstrated inside a hollow-core photonic crystal fiber in the first part of this thesis. However, the coherence time of the superposition state is limited by the inhomogeneous broadening from the optical trapping potential to tens of microseconds [52]. Here, we studied the coherence properties of ^{85}Rb atoms confined and guided inside a hollow-core photonic crystal fiber.

By cancelling the differential ac stark shift from the optical trapping potential, we first guided a quantum superposition state of the trapped atoms inside the fiber over a centimeter distance and preserved the coherence over hundreds of milliseconds. The coherence time and transportation distance are enhanced by orders of magnitude compare to the previous schemes [44].

1.4 Overview

This thesis presented as following. Chapter 2 provides a brief introduction of atom-light interaction, atom cooling and trapping techniques, and detection of atom number and temperature. Chapter 3 gives an introduction of atomic sensors based on matter-wave interferometer, and the gravity measurement by using matter-wave interferometer. The path integral approach is used to discuss the phase shifts in the interferometer. In Chapter 4, I introduce the experimental apparatus used for the experiments, including fiber selection, vacuum system, optical setup, magneto-optical trap, control systems, and sensor hardware used in this work. Chapter 5 describes the process of loading cold atoms into a hollow-core fiber, and the results illustrate that atoms are trapped into the fiber. Chapter 6 provides the results and discussions of an atom interferometer inside the hollow-core fiber. In Chapter 7, transporting long-lived quantum spin coherence inside a fiber is presented. Finally, Chapter 8 summarizes the results of our work, and describes general outlooks to improve the performance of atom interferometers inside the hollow-core fiber. Chapter 9 is the appendix, including characteristic data of ^{85}Rb , fiber preparation, design of the magnetic coils and water cooling system, and derivation of the number and temperature of atoms inside fiber.

Chapter 2

Laser cooling and trapping

This chapter starts from the description of atomic structure of ^{85}Rb followed by the interaction between a two-level atom and light. Understanding the mechanisms of atom-light interaction in two-level atomic system provides a significant conceptual framework for atom cooling and trapping. The optical forces used to cool and trap atoms are described. Then, the techniques of magneto-optical trap (MOT) are detailed presented [53, 54]. Finally, the methods of number and temperature measurements are given. These techniques allow us to cool neutral rubidium atoms from room-temperature to $\sim 10\ \mu\text{K}$ cold atomic cloud, and trap them into a hollow-core fiber.

2.1 Atomic structure

We used alkali-metal atoms, such as the ^{85}Rb atoms are used in our experiments, the atomic structure consists of a closed inner shell and one outermost electron in the S-orbital. The simple energy levels allow for clean and efficient atom-light interactions. They can be cooled efficiently with established laser cooling technology. Rubidium is one of the most popular atomic species for laser cooling and trapping, and ^{85}Rb isotope has the largest natural abundance Rubidium which is beneficial for trapping large atom number. It can also be easily converted to or combine with ^{87}Rb for other interesting experiments, such as testing universality of free fall of atoms.

Figure 2.1: Rubidium (^{85}Rb) atom energy level structure (The D₁ line transition ($5^2S_{1/2} \rightarrow 5^2P_{1/2}$) and the D₂ line transition ($5^2S_{1/2} \rightarrow 5^2P_{3/2}$)).

The simplified energy level splitting diagram of ^{85}Rb is shown in figure 2.1. It is a fine-structure doublet consisting of the $5^2S_{1/2} \rightarrow 5^2P_{3/2}$ and $5^2S_{1/2} \rightarrow 5^2P_{1/2}$ transitions. The fine-structure doublet is a result of the total electron angular momentum $\mathbf{J} = \mathbf{L} + \mathbf{S}$, where $\mathbf{S}$ is the spin angular momentum and $\mathbf{L}$ is the orbital angular momentum of the outermost electron. For the ground state of ^{85}Rb , where $S = 1/2$ and $L = 0$, the total electron angular momentum $J = 1/2$. For the first excited states, the orbital angular momentum $L = 1$, so $J = 1/2$ or $3/2$. Therefore, the transition from $L = 0$ to $L = 1$ is split into two components, the D₁ line transition ($5^2S_{1/2} \rightarrow 5^2P_{1/2}$) and the D₂ line transition ($5^2S_{1/2} \rightarrow 5^2P_{3/2}$).

For the hyperfine structure of ^{85}Rb , it is a result of the total atomic angular momentum $\mathbf{F} = \mathbf{J} + \mathbf{I}$. It is the coupling of the total electron angular momentum $\mathbf{J}$ with the nuclear angular momentum $\mathbf{I}$. The ^{85}Rb atom has a nuclear spin $I = 5/2$. For the ground state, $J = 1/2$, the total angular momentum F is 2 or 3. The ground state is split into two states according to different values of F . For the excited state of D₁ line ($5^2\text{P}_{1/2}$), $J = 1/2$, so the total angular momentum F' is 2 or 3. For the D₂ excited state ($5^2\text{P}_{3/2}$), $J = 3/2$, so F can take any of the values 1, 2, 3, or 4. The D₁ line transition ($5^2\text{S}_{1/2} \rightarrow 5^2\text{P}_{1/2}$) can be driven by laser around wavelength $\lambda \approx 795$ nm and the D₂ line transition ($5^2\text{S}_{1/2} \rightarrow 5^2\text{P}_{3/2}$) can be driven by laser around wavelength $\lambda \approx 780$ nm.

Each hyperfine state can be split into $2F + 1$ Zeeman sub-levels by applying an external magnetic field, which is called Zeeman effect. The so-called “clock transition” between the $m_F = 0$ states of the two ground states $F = 2$ and $F = 3$ ($|F = 2, m_F = 0\rangle \rightarrow |F = 3, m_F = 0\rangle$) is usually used in atom interferometer, as the clock states are only sensitive to second order Zeeman effect. More details about the atomic structure and the physical and optical properties of ^{85}Rb are presented in Appendix A.

2.2 Atom-light interaction

We start by introducing some details regarding atom-light interactions before the description of laser cooling and trapping. We first consider a single atom with two states (Fig. 2.2) interacting with a single monochromatic electric field. The ground state is $|g\rangle$, the excited state is $|e\rangle$, and the energy separation is $\hbar\omega_{eg}$, where ω_{eg} is the resonance frequency which drives the transition between state $|g\rangle$ and state $|e\rangle$. The frequency of the electric field is ω_L . Consider a two-level atom of the unperturbed Hamiltonian:

$$H_0 = \hbar \begin{pmatrix} \omega_e & 0 \\ 0 & \omega_g \end{pmatrix}. \quad (2.1)$$

Figure 2.2: Interaction between a two-level atom and a monochromatic electric field. The energy difference between the ground and excited states is $\hbar\omega_{eg}$, and Γ is the decay rate of the excited state.

Using the semi-classical approximation method, the Hamiltonian for the two-level atom in the presence of the monochromatic electric field is given by:

$$H = H_0 + H' = H_0 + \mathbf{d} \cdot \mathbf{E} \quad (2.2)$$

where $\mathbf{d}$ is the dipole momentum operator and the monochromatic electric field is

$$\mathbf{E} = E_0 \hat{e} \cos(\mathbf{k} \cdot \mathbf{r} + \omega_L t + \phi). \quad (2.3)$$

The interaction between the atom and the electric field is given as an off-diagonal matrix element of:

$$\Omega = \frac{-\langle e | \mathbf{d} \cdot \mathbf{E} | g \rangle}{\hbar} \quad (2.4)$$

which is the Rabi frequency of the oscillation between the $|g\rangle$ state and $|e\rangle$ state for the condition $\omega_L = \omega_{eg} = \omega_e - \omega_g$.

Using the Hamiltonian, the time-dependent Schrödinger equation leads to:

$$\dot{c}_g = -i \frac{\Omega}{2} (e^{-i(\omega_{eg} - \omega_L)t} + e^{-i(\omega_{eg} + \omega_L)t}) c_e \quad (2.5)$$

$$\dot{c}_e = -i\frac{\Omega}{2}(e^{i(\omega_{eg}-\omega_L)t} + e^{i(\omega_{eg}+\omega_L)t})c_g \quad (2.6)$$

where c_g and c_e are the coefficients of the state vector in two-level atom. By considering the condition $\omega_L = \omega_{eg}$, and using the rotating wave approximation, the equations of time evolution are now:

$$\dot{c}_g = -i\frac{\Omega}{2}e^{-i(\omega_{eg}-\omega_L)t}c_e \quad (2.7)$$

$$\dot{c}_e = -i\frac{\Omega}{2}e^{i(\omega_{eg}-\omega_L)t}c_g. \quad (2.8)$$

Choosing the initial conditions that atoms are in ground state $|g\rangle$ initially, i.e. $c_g(0) = 1$ and $c_e(0) = 0$, we get:

$$c_e(t) = -i\sin\left(\frac{\Omega}{2}t\right). \quad (2.9)$$

Therefore, the probability of being in the $|e\rangle$ state (i.e. the upper state) is therefore

$$|c_e(t)|^2 = \sin^2\left(\frac{\Omega}{2}t\right). \quad (2.10)$$

The two-level system is oscillating at a period of $2\pi/\Omega$. This phenomena is called Rabi oscillation. If we consider the condition that the light is off-resonance ($\omega_L \neq \omega_{eg}$) and still with the initial condition: $c_g(0) = 1$ and $c_e(0) = 0$. The probability of being in the $|e\rangle$ state is

$$|c_e(t)|^2 = \frac{\Omega^2}{\Omega'^2} \sin^2\left(\frac{\Omega'}{2}t\right) \quad (2.11)$$

where Ω' is the generalized Rabi frequency of the Rabi oscillation. We define the detuning from resonance as $\Delta = \omega_L - \omega_{eg}$, the Rabi frequency is in the general expression $\Omega'^2 = \Omega^2 + \Delta^2$. We also observe an oscillation among the populations of the two-level system. Compared to the resonant case, the transfer to the excited state is not completed during the oscillation. We can get a similar formula of the probability to find atom in the ground state if atom

is in the excited state $|e\rangle$ initially. At the resonance case $\omega_L = \omega_{eg}$, an atom in the ground or excited state interacting with the light for a time $t = \pi/\Omega$ will be fully transferred to the opposite state. The pulse $t = \pi/\Omega$ is typically called a π pulse. If the atom-light interaction last for a time $t = \pi/2\Omega$, the atom will be transferred to the other state with 50% probability; such a pulse is typically called a $\pi/2$ pulse. In practice, the light will introduce dephasing, such as spontaneous emission and differential ac Stark shift, during the interaction, which causes the amplitude of the Rabi oscillation to decrease. The dephasing mechanism will be detailed discussed in chapter 7.

For a time $t \ll 1/\Omega'$, the atoms are mainly in the ground state $|g\rangle$. In this condition we can define the absorption rate Γ_A as the rate the system leave the ground state $|g\rangle$ due to the interaction with the electromagnetic field.

$$\Gamma_A = \frac{d}{dt} |c_e(t \ll 1/\Omega')|^2 = \frac{\Omega'^2}{2} t. \quad (2.12)$$

2.3 Optical forces

The force acting on atoms due to the atom-light interaction has two terms: scattering force (or radiative force) and dipole force [54]. The first term is from the absorption and spontaneous emission process, and the second term is a conservative component due to photon redistribution among the plane waves. In this section, this two primary forces are detailed discussed.

2.3.1 Scattering force

We first again consider a two-level atom interacting with light. If the frequency of the light is close to the resonant frequency ω_{eg} ($\omega_L \approx \omega_{eg}$), and the atom is initially in ground state $|g\rangle$, the atom will absorb a photon and go to excited state $|e\rangle$. Once it is in excited state $|e\rangle$, the atom can spontaneously emit a photon into an electrical field of different mode with the frequency ω_S , which occurs at a rate of $\Gamma/2$, where Γ is the linewidth of the atomic resonance transition. According to momentum conservation, when a photon is absorbed, its momentum $\Delta\mathbf{p} = \hbar\mathbf{k}$ is transferred to the atom retaining both its magni-

Figure 2.3: (a) Simplified diagram for describing scattering force on a two-level atom. (b) The schematics of spontaneous emission of an atom (atom will emit photons in random direction).

tude and direction, where $\mathbf{k}$ is the wave vector; $k = 2\pi/\lambda$ is the wave number of the optical field; λ is the wavelength) (see Fig. 2.3). Since the spontaneous emission is isotropic (atom emits photons into random directions), the momentum recoil of the atom related to the spontaneous emission is in an arbitrary direction. On average, the change of momentum is zero. That means, if an absorption is followed by a spontaneous emission (the wave vector of the spontaneous emission light is $\mathbf{k}_s$), the atom gains momentum

$$\Delta\mathbf{p} = \hbar\mathbf{k} - \hbar\mathbf{k}_s \quad (2.13)$$

With an expectation value of

$$\langle\Delta\mathbf{p}\rangle = \hbar\mathbf{k} - \hbar\langle\mathbf{k}_s\rangle = \hbar\mathbf{k}. \quad (2.14)$$

Thus the overall change of the atomic momentum is solely from the effect of absorption (see Fig. 2.3). After n times of absorption-spontaneous emission cycles, the atom gains a momentum $Mv_{rec} = n\hbar k$ along the propagation direction of incoming photons, where M is the atom mass and v_{rec} is the recoil velocity from photons. For a Rubidium atom, the recoil velocity v_{rec} at 780 nm wavelength is about 3.5 mm/s, which is much smaller than the thermal

velocity of atoms at room temperature. In such a condition an atom moving in the direction opposite to the beam of the photon would be slowed down by the transfer of the momentum during the absorption-emission process. This recoil force acting on the atom is the scattering force. The scattering rate is given by:

$$\Gamma_{scat} = \frac{\Gamma}{2} \cdot \frac{I/I_{sat}}{1 + I/I_{sat} + 4(\Delta/\Gamma)^2} \quad (2.15)$$

and the scattering force is

$$\mathbf{F} = \hbar \mathbf{k} \Gamma_{scat} = \frac{\hbar \mathbf{k} \Gamma}{2} \cdot \frac{I/I_{sat}}{1 + I/I_{sat} + 4(\Delta/\Gamma)^2} \quad (2.16)$$

where I is the intensity of the light, Δ is the detuning from resonance ($\Delta = \omega_L - \omega_{eg}$), and I_{sat} is the saturation intensity of the atomic resonance transition. Absorptions followed by spontaneous emissions are at the origin of the scattering force induced by atom-light interaction. The scattering force is a dissipative force since atoms couple to the vacuum field through spontaneous emission.

2.3.2 Dipole force

Optical dipole force is from the electric dipole of atoms interacting with an off-resonance light. The energy levels of the atomic states are shifted due to the interaction with the light field, also known as ac Stark shift [55].

$$\Delta E = \pm \frac{3\pi c^2}{2\omega_{eg}^3} \frac{\Gamma}{\Delta} I(\mathbf{r}) \quad (2.17)$$

where $I = I(\mathbf{r})$ is the spatially dependent light intensity. This term is proportional to the spatial profile of the laser intensity. For red-detuned light fields, $\Delta < 0$, atoms are pushed toward the high intensity region of the light, while for blue-detuned light fields, $\Delta > 0$, atoms are pushed toward the low intensity areas of the light (see Fig. 2.4).

We consider a single beam with a Gaussian beam profile and the detuning of the light $\Delta < 0$. This Gaussian beam can create a harmonic potential for

Figure 2.4: Dipole potentials and energy shifts of a dipole trap. For red detuning ($\Delta < 0$), an attractive force for the atoms can be created from a Gaussian beam. For a blue detuning ($\Delta > 0$), a hollow beam creates an attractive force at the center of the beam for atoms.

atoms in the high intensity areas (center region). The expression of the optical dipole potential is given by

$$U_{dip}(\mathbf{r}) = \frac{3\pi c^2}{2\omega_{eg}^3} \frac{\Gamma}{\Delta} I_0 \quad (2.18)$$

where $I = I_0$ is the maximum intensity of the Gaussian beam. Atoms will experience a force given by the gradient of the optical dipole potential introduced by the spatially varying trapping beam. Optical dipole traps are usually used to load cold atoms from a MOT which provides a way to trap cold atoms after cooling with low single-photon scattering.

For optical dipole traps in free space, atoms' lifetime in the trap is mainly limited by the scattering from the optical trapping beam. The scattering rate from the dipole trapping beam is

$$\Gamma_{sc} = \frac{3\pi c^2}{2\omega_{eg}^3} \left(\frac{\Gamma}{\Delta}\right)^2 I(\mathbf{r}). \quad (2.19)$$

The key parameters of optical dipole trapping for far-detuned light fields can be understood by these two equations (Eq. 2.18, 2.19). The dipole potential scales as $1/\Delta$ and the scattering rate is proportional $\sim 1/\Delta^2$. Therefore, the heating from the scattering can be reducing to a negligible level by increasing

the detuning of the trapping light frequency. Compared to red-detuned dipole traps with Gaussian beam profiles, blue-detuned dipole traps can also be implemented with donut shape beam profiles, where the intensity minimum is at the center of the beam. The main advantage of the blue-detuned traps is that atoms mostly stay in the dark region of the trapping beams and thus the heating introduced by scattering can be significantly reduced.

The AC Stark shift varies as a function of the spatial distribution of the intensity. Therefore, the inhomogeneous intensity distribution of the trapping light and the distribution of atoms in the trapping potential introduce inhomogeneous broadening on the transition when the two states have different ac Stark shifts. This dephasing mechanism is the main limitation to spin-coherence times. The detailed discussion of the differential ac Stark and the cancellation of the differential shift is presented in chapter 7.

2.4 Laser cooling

2.4.1 Doppler cooling and optical molasses

A briefing introduction on the mechanism of Doppler cooling and optical molasses is presented here. To cool atoms, we should consider the scattering rate corresponding to the velocity of atoms by using the Doppler effect. Let us imagine this situation. If a two-level atom of velocity v propagates against ($\mathbf{k} \cdot \mathbf{v} < 0$) a photon beam with the frequency ω_L , and the frequency of photons is slightly tuned below the transition resonance of frequency ω_{eg} ($\omega_L < \omega_{eg}$), the traveling atoms see the frequency related to the Doppler shift

$$\omega = \omega_L - \mathbf{k} \cdot \mathbf{v} \quad (2.20)$$

which is closer to atomic resonance. Therefore, faster atoms will have higher probability to absorb photons to decelerate.

On the contrary, if atoms propagates along ($\mathbf{k} \cdot \mathbf{v} > 0$) the propagating direction of the photon beam, the detuning from resonant transition is even further increased because of the Doppler effect ($\omega = \omega_L - \mathbf{k} \cdot \mathbf{v}$) and the

Figure 2.5: A schematic of the principle of Doppler cooling and optical molasses. The laser frequency, which is tuned slightly below the resonant transition, is shifted due to the velocity of the atom.

probability of the absorption for faster atoms is reduced. The photon detuning depending on the direction of its motion results in an asymmetric influence, and atoms are slowed down by the deceleration and acceleration forces from the two counter-propagating photon beams (see Fig. 2.5). We can extend the slowing effect to three dimensions by using three intersecting orthogonal pairs of counter-propagating laser lights [56]. Each atom moving in this overlapping region of beams will be slowed down. The slowing effect will vanish only if atoms move out of the region.

The total optical pressure force $\mathbf{F}$ from the two beams is given by

$$\mathbf{F} = \mathbf{F}_+ + \mathbf{F}_- \quad (2.21)$$

with

$$\mathbf{F}_{\pm} = \pm \frac{\hbar \mathbf{k} \Gamma}{2} \frac{S_0}{1 + S_0 + [2(\Delta \mp \mathbf{k} \cdot \mathbf{v})/\Gamma]^2}. \quad (2.22)$$

The total force acting on the atom is

$$\mathbf{F} = \frac{8\hbar k^2 \Delta S_0 \mathbf{v}}{\Gamma[1 + S_0 + (2\Delta/\Gamma)^2]^2} \quad (2.23)$$

where $S_0 = I/I_{sat}$, is the saturation parameter, I is the light intensity, $I_{sat} = \pi \hbar c \Gamma / 3 \lambda^3$ is the saturation intensity of the resonance transition, Γ is the decay

Figure 2.6: The optical forces versus velocity for one dimensional optical molasses (for $\Delta < 0$). Dashed lines show the two component forces $\mathbf{F}_+$ (red line) and $\mathbf{F}_-$ (blue line) while the solid line shows the total force acting on atoms.

rate, Δ is the detuning from the resonance. Figure 2.6 shows the relationship of the force acting on the atoms and the velocity of atoms. For $\Delta < 0$, when the velocity is positive, the total force is negative; when the velocity is negative, the total force is positive. The force given by the laser beams always slows down the motion of atoms. This is called optical molasses (OM) [57]. Doppler cooling can happen only in the condition $\Delta < 0$, in the opposite condition (for $\Delta > 0$) the atom will heat up due to the acceleration. We can find a region where the force is linearly changed with the velocity of atoms. The total force $\mathbf{F}$ is proportional to velocity which is given by an approximating expression [58]:

$$\mathbf{F} = \frac{8\hbar k^2 \Delta S_0 \mathbf{v}}{\Gamma[1 + S_0 + (2\Delta/\Gamma)^2]^2} \approx -\beta \mathbf{v} \quad (2.24)$$

if assuming the velocity of atoms satisfy the condition $(kv/\Gamma)^4 \ll 1$. This type of pressure force given by Eq. (2.24) alone cannot achieve a stable trapping of

atoms [59]. The photon recoils from the scattering of the laser beams results in that atoms would move randomly in the optical molasses. This random walk due to the scattering limits the temperature achieved by Doppler cooling to be

$$T_D = \frac{\hbar\Gamma}{2k_B}. \quad (2.25)$$

Additionally, there is only viscous confinement force in the optical molasses. Atoms will eventually escape from the optical molasses due to the random motion. That is why the Doppler cooling only cannot confine atoms.

2.4.2 Sisyphus cooling

The temperature of atoms to be much lower than the Doppler cooling limit T_D in the optical molasses [60] can actually be achieved by an additional cooling method known as Sisyphus cooling or polarization gradient cooling. The cooling mechanism of Sisyphus cooling fundamentally differs from the Doppler cooling. If the polarization of the light field is spatially varying, atoms in the optical molasses would move in a non-adiabatic path. Such a spatially varying light field can be created by two counter-propagating laser beams with right circular σ^+ and left circular σ^- polarizations ($\sigma^+ - \sigma^-$).

The net polarization of the two counter-propagating beams ($\sigma^+ - \sigma^-$) is linear with spatially rotating in a period of $\lambda/4$. The working principle of Sisyphus cooling is based on the ac Stark shift introduced by the two counter-propagating laser beams. The spatial rotation of the polarization leads to an ac Stark shift varying in space at a period of $\lambda/4$. Figure 2.8 shows a simplified description of the mechanism. Atoms moving through the spatially modulated potential along the optical lattice will lose their kinetic energy and change their potential energy levels by interacting with laser beams. When atoms travel from the “valley” to the “hill”, they will lose their kinetic energy. As atoms arrive at the “hill”, the optical pumping beam will pump atoms to the “valley” and atoms also lose their potential energy. In this way, atoms convert their kinetic energy to the potential energy, and lose their potential energy by interacting with laser beams. This method has been used to cool atoms below

Figure 2.7: A specified schematic to explain the mechanism of the Sisyphus cooling. Atoms moving through the spatially modulated potential along the optical lattice will lose their kinetic energy and change their energy levels by interacting with laser beams.

the Doppler cooling [61, 62].

The temperature can be reached by Sisyphus cooling scales as

$$T_S = \frac{\hbar \Omega'^2}{k_B \Delta} \quad (2.26)$$

where Ω' is the Rabi frequency, which is mentioned in section 2.2, Δ is the detuning frequency to atomic resonance transition, and k_B is the Boltzmann constant. This cooling method also has a limit. Temperature can be achieved by Sisyphus cooling is limited by the photon recoil,

$$T_{rec} = \frac{(\hbar k)^2}{2k_B m} \quad (2.27)$$

where k is the wave vector of the laser beam and m is the mass of atom. Similar to Doppler cooling, there is no only restoring confinement in the Sisyphus cooling but only viscous force. Therefore atoms can eventually escape the optical molasses region due to the random motion.

2.5 Magneto-optical trap (MOT)

2.5.1 Trap loading

Doppler cooling and Sisyphus cooling are effective methods to cool atomic cloud to a low temperature. However, they are not able to provide confinement of atoms in a well-defined region. It's easy to understand: for any trapping volume in an optical molasses, the optical forces in this region is the same, thus there is no position dependent optical forces pointing towards to the center of this region.

The simultaneous cooling and confinement of atoms can be achieved by the magneto-optical trap which uses a magnetic field to produce Zeeman shifts of the energy levels of atoms. To understand the mechanism of a magneto-optical trap, we consider a simple two-level system in one-dimensional case, illustrated in Fig. 2.8. It consists of two red-detuned counter-propagating beams with opposite circular polarization (right σ^+ and left σ^-). The magnetic field

Figure 2.8: (a) Arrangement of a one-dimensional magneto-optical trap. (b) Atoms in an external magnetic field interact with laser beams. Atoms are trapped in the field center where the magnetic field is zero. $F = 0$ is the ground state and $F' = 1$ is the excited state.

$\mathbf{B}$ linearly in magnitude along z axis. In Fig. 2.5(b), due to the selection rules, the σ^+ (or σ^-) circularly polarized beam causes only transitions with $\Delta m_F = +1$ (or $\Delta m_F = -1$). Under the influence of magnetic field, the excited state $F' = 1$ is split to three Zeeman sub-levels (known as Zeeman effect). The energy shift of the sub-levels is related to its magnetic quantum number m_F and the magnetic field B :

$$\Delta E_{m_F} = g_F \mu_B B m_F, (m_F = 0, \pm 1, \dots, \pm F) \quad (2.28)$$

where g_F is Landé factor and μ_B is Bohr magneton.

Firstly, we consider atoms moving in the left region ($-z$ region). The beam from the left is right circular polarized (σ^+). The Doppler effect will modify the frequency atoms see, so the σ^+ beam is tuned closer to the resonance transition. The beam sent from right side with σ^- polarization is even further detuned from the resonance by Doppler effect in the left region. Therefore the σ^+ polarized beam mainly excites the transitions between the states $F = 0$ and $F' = 1$. Due to the momentum transferred from σ^+ beam to the atom, the

Figure 2.9: A three-dimensional magneto-optical trap. A pair of anti-Helmholtz coils generates a spherical-quadrupole-symmetry magnetic field in all three directions.

atom is pushed toward the center ($B = 0$). For the right region, the momentum transfer of σ^- will also push atom to the center region. As a result, atoms are always pushed toward the trap center. The Doppler and sisyphus cooling are responsible for cooling atoms and Zeeman effect provides spatially dependent force for trapping. The magneto-optical trap described above can be extended to the three-dimension by simply adding another two pairs of laser beams in orthogonal directions as shown in Fig. 2.9.

2.5.2 Number of atoms

There are two key parameters of atoms in the MOT that are most interested for the experiments, the number (or atomic density) and temperature of atoms. In this section, I will present the method used to measure the number of atoms in our experiment.

Figure 2.10: A fluorescence detection system used to measure the number of atoms. The fluorescence signal of atoms is detected by a sensitive photodetector. The fluorescence photons are focused onto the photodetector by using a lens with diameter $D = 50$ mm, and the distance L between the lens and the atomic cloud is 150 mm.

There are two common methods used to measure the number of atoms. One is detecting the fluorescence signal scattered from the trapped atoms in the MOT where the scattered light is proportional to the number of atoms. The other one is measuring the absorption signal of a probe beam that passes through the atoms.

In this experiment, we choose the fluorescence method to measure the number of atoms outside the fiber. The fluorescence photons are focused to a CCD by using a lens of diameter $D = 50$ mm and focal length $f = 100$ mm. The distance L between the lens and the center of the atomic cloud is 150 mm. The detection system is shown in Fig. 2.10. By converting the response of the CCD from voltage to optical power, we can determine the optical power of the fluorescence photons.

To find the total number of atoms, we must obtain the total optical power emitted by the atomic cloud. Therefore, we first calculate the ratio of the optical power detected and the total power emitted which can be found by

the ratio between the solid angle of the focusing lens and the solid angle of a sphere ($\Omega(\theta = \pi) = 4\pi$). The solid angle of the detection system is $\Omega(\theta) = 2\pi(1 - \cos \theta)$, where $\theta = \arctan(\frac{D}{2L})$. Hence, the ratio is

$$\frac{\Omega(\theta)}{\Omega(\theta = \pi)} = \frac{1 - \cos \theta}{2}. \quad (2.29)$$

The energy of one photon is $E = \hbar\omega = hc/\lambda$, and the scattering rate is given in Eq. 2.15. In this case, I is the total intensity of the trapping beams, and I_{sat} is the saturation intensity of the tapping beams. The total number of atoms in the atomic cloud is calculated from the data is given by

$$N_{atom} = \frac{2}{(1 - \cos \theta)} \frac{1}{\Gamma_{scat}} \frac{P}{h\lambda/c} \quad (2.30)$$

where P is the detected optical power of the fluorescence photons in the detection system.

Once the number of atoms is detected, we can estimate the atomic density of the atomic cloud. We assume the ideal case, the intensity of all cooling beams are equal, and the gradient of magnetic field on the z axis is twice that of the gradient on the $x - y$ plane. In this ideal case, the shape of this atomic cloud is like an ellipsoid with $r_x = r_y = \sqrt{2}r_z$ (r_x , r_y and r_z is the radius of the atomic cloud in three dimensions respectively). Hence, the atomic density of the cloud is

$$n = \frac{3N_{atom}}{\pi(2r_z)^3}. \quad (2.31)$$

If we assume the atomic cloud has a spherical symmetry, the expression of the atomic density can be further simplified as

$$n = \frac{3N_{atom}}{4\pi r_z^3}. \quad (2.32)$$

2.5.3 Temperature measurement

Figure 2.11: The time-of-flight method for measuring the temperature of trapped atoms. The fluorescence images of the expanding atomic cloud are collected by a CCD camera, and the density distribution is extracted from these images. The MOT beams are used to excite the fluorescence of atoms. The time-of-flight method for measuring the temperature of trapped atoms.

We use the time-of-flight (TOF) method to measure the temperature of cold trapped atoms [63]. This method measures the density distribution of expanding atomic cloud at different time in free space when it is released from the trap. A simplified schematic illustrates the time-of-flight detection in Fig. 2.11. The distribution of the density with temperature T during free falling along y direction is given by [64]

$$\rho(x, y, t) = \frac{1}{2\pi\sigma^2(t)} \exp\left(-\frac{y^2(t) + x^2}{2\sigma^2(t)}\right) \quad (2.33)$$

where $\sigma(t)$ is the standard deviation of the density distribution, and the center position of atomic cloud in y direction is

$$y(t) = y_0 + \frac{1}{2}gt^2 \quad (2.34)$$

with y_0 is the initial position in y direction. The time evolution of the deviation is determined by

$$\sigma^2(t) = \sigma_0^2 + (k_B T/m)t^2 \quad (2.35)$$

with σ_0 is the standard deviation of the density distribution at time $t = 0$. This equation determines the expansion of the atomic cloud due to its temperature T . Once we obtain the evolution of the density distribution, we can calculate the temperature of the atoms according to Eq. 2.35.

Chapter 3

Atom interferometry

Matter wave interferometers using ultra-cold atoms have evolved into a mature tool for precision measurement and quantum sensing since the first demonstration in 1991 [65, 66]. Atom interferometers use both wave and particle properties of atom to perform measurements. In this chapter, I will give an overview of atom interferometry and introduce its application for precision measurement of gravity g .

3.1 Ramsey interferometer

In section 2.2, we discuss the interaction between a two-level atom and a monochromatic electromagnetic (EM) field. In the presence of EM field with Rabi frequency Ω , the two-level system is oscillating at a period of $2\pi/\Omega$. On resonance, an atom in the ground or excited state interacting with the EM field for a time $t = \pi/\Omega$ will be fully transferred to the other state with unity probability. The pulse $t = \pi/\Omega$ is typically called a π – pulse. If the atom-light interaction last for a time $t = \pi/2\Omega$, the atom will be transferred to the other state with 50% probability, such a pulse is typically called a $\pi/2$ – pulse. If atoms are atom initially in the ground state $|g\rangle$. The π – pulse and $\pi/2$ – pulse corresponds to the operation

$$t = \pi/\Omega, \pi - \text{pulse}, |g\rangle \rightarrow |e\rangle$$

Figure 3.1: (a) The time sequence of Ramsey interferometer. A $\pi/2 - \text{pulse}$ is used to prepare the superposition state, then a second $\pi/2 - \text{pulse}$ is applied after allowing atoms to evolve for a separation time T . (b) Schematic of a Ramsey interferometer using microwave pulses.

$$t = \pi/2\Omega, \pi/2 - \text{pulse}, |g\rangle \rightarrow (|g\rangle + |e\rangle)/\sqrt{2}. \quad (3.1)$$

A Ramsey interferometer makes use of microwave instead of light to interact with atoms. It was first demonstrated by Norman Ramsey in 1950 [67,68] who was awarded the Nobel Prize in Physics in 1989.

In a microwave Ramsey interferometer, atoms are prepared in the ground state $|g\rangle$, and a $\pi/2 - \text{pulse}$ is applied to put the atoms into a coherent superposition of ground and excited state. The phases of atoms in different energy states will evolve in time at different rates. After a time T , a second $\pi/2 - \text{pulse}$ is applied to bring atoms to excited state and interfere. If there is no detuning, total atoms would go to the excited state, which can be detected. If there is no detuning between the microwave frequency and the transition frequency, all atoms will go to the excited state for detection. If there is a finite detuning, the population of atoms in the excited will be determined by the phase $\varphi = \Delta \cdot T$.

The interference fringes can be obtained by measuring the population in the excited or ground state as a function of the separation time T , detuning Δ or the relative phase between the two $\pi/2 - \text{pulses}$. If there is a dephasing effect during the separation time, the contrast of the interference fringes would decay as a function of T . The decay constant is usually used to define the coherence of the atomic system. In this experiment, we used the Ramsey method to

measure the coherence time of atoms and study its dephasing mechanism in the optical dipole trap formed by the fundamental mode of the fiber.

3.2 Light-pulse atom interferometry

Instead of using microwave pulses in Ramsey method, the optical pulses are applied in Doppler sensitive fashion to place atoms into a superposition state, and split and recombine atoms spatially. Opposed to Ramsey interferometers using microwave, atomic states in light-pulse interferometers have non-negligible spatial separation. In this section, I will discuss the working principle of this type of interferometers and calculate the phase shift in Mach-Zehnder geometry [69].

3.2.1 Raman Mach-Zehnder interferometer

Figure 3.2: Space-time diagram of a Mach-Zehnder interferometer based on three Raman pulses. There is no external forces acting on the atoms. The first Raman pulses is applied to split atoms to two trajectories followed by two more pulses to reflect and recombine them for interference.

Doppler sensitive light-pulse atom interferometers use two-photon stimulated Raman transitions [70, 71] to address a pair of long-lived states. In Mach-Zehnder atom interferometers, the three Raman pulses are applied to split atoms to two trajectories, and reflect and recombine the atomic wave-packets to interfere. Between the Raman pulses, the atoms are free-falling in order to increase the interrogation time.

In light interferometers, light is split, recombined, and interfered through two spatially distinct trajectories. Similar to light interferometers, matter waves are initially split to two spatially separated trajectories by a beam splitter pulse. Atoms in the two arms are then reflected by a mirror pulse followed by another beam splitter pulse for recombination to form interference (see Fig. 3.2). The probability of finding atoms in the state $|1\rangle$ or state $|2\rangle$ depends on the phases between the two interferometer arms. For our experiments, we used ^{85}Rb to perform the Mach-Zehnder atom interferometers, state $|1\rangle$ or state $|2\rangle$ are the hyperfine ground states state $|F = 2\rangle$ or state $|F = 3\rangle$ respectively.

The principle of two-photon stimulated Raman transitions is illustrated in Fig. 3.3. Consider the first beamsplitter at $t = 0$ in Fig. 3.2, atoms are initially in state $|1\rangle$ with momentum $\mathbf{p}$. Atoms absorb the laser light sent from the left with wave-vector $\mathbf{k}_1$ will receive a momentum kick $\hbar\mathbf{k}_1$ from the light. Atoms will then stimulated emit a photon due to the second laser beam with wave-vector $\mathbf{k}_2$, sent from the right and receive a momentum kick $\hbar\mathbf{k}_2$. During this procedure, the momentum of atoms increase with the change of internal state. This two-photon process Raman process brings atoms to state $|2\rangle$ with a recoil momentum kick

$$\Delta\mathbf{p} = \hbar\mathbf{k}_1 - \hbar\mathbf{k}_2 = \hbar\mathbf{k}_{eff} \approx 2\hbar\mathbf{k}_1. \quad (3.2)$$

In the two-photon Raman transitions, the two hyperfine ground states $|1\rangle$ and $|2\rangle$ are driven by the two counter-propagating beams through a virtual state (the detuning to the excited state is Δ). We can consider this system as an effective two-level system for simplification. Choosing the initial conditions that atoms are in ground state $|1\rangle$ initially. The probability of being in the

Figure 3.3: Stimulated Raman transition process of a three-level atom. (a) A two-level system with two hyperfine ground states $|1\rangle$ and $|2\rangle$. (b) Two-photon stimulated Raman transition process. Two counter-propagating laser beams with wave-vector $\mathbf{k}_1$ and wave-vector $\mathbf{k}_2$, interacting with atoms by two-photon Raman process.

two states after interacting with a Raman pulse are

$$|c_1(t)|^2 = \cos^2\left(\frac{\Omega_{eff}}{2}t\right) \quad (3.3)$$

$$|c_2(t)|^2 = \sin^2\left(\frac{\Omega_{eff}}{2}t\right) \quad (3.4)$$

where Ω_{eff} is the two-photon Rabi frequency, which is given by

$$\Omega_{eff} = \frac{\Omega_1 \Omega_2}{2\Delta} \quad (3.5)$$

where Ω_1 and Ω_2 is the Rabi frequencies of the two laser beams respectively, Δ is the detuning from the excited state.

For an optical pulse with duration $t = \pi/2\Omega$, $|c_1(t)|^2 = |c_2(t)|^2 = 1/2$ and atoms are transferred to the state $|2, \mathbf{p} + \hbar\mathbf{k}\rangle$, with 50% probability. Such a pulse is the $\pi/2$ -pulse and acts as a beamsplitter. For a pulse with duration $t = \pi/\Omega$, atoms in the ground state $|1, \mathbf{p}\rangle$ interacting with the two laser beams are fully transferred to the state $|2, \mathbf{p} + \hbar\mathbf{k}\rangle$. The pulse $t = \pi/\Omega$ is the π -pulse and acts as a mirror. Figure 3.2 shows how the three pulses ($\pi/2 - \pi - \pi/2$) is arranged in a Mach-Zehnder atom interferometer. The first $\pi/2$ -pulse works

as a beamsplitter to split the wavepackets into two paths. The π – pulse is used to reflect and recombine atoms in different trajectories. Finally, the last $\pi/2$ – pulse is sent to interfere atoms.

After the second $\pi/2$ – pulse, atoms in the two output arms of the interferometers have different momentum $\mathbf{p}$ and $\mathbf{p} + \hbar\mathbf{k}$. Therefore, this two branches of atoms will be separated in space over time and can be detected by fluorescence imaging or absorption imaging. In the following section, I will discuss the phase in more details.

3.2.2 Interferometer phase shifts

In a Mach-Zehnder type atom interferometer, after a $\pi/2$ – pulse, a π – pulse and a second $\pi/2$ – pulse at time t_1 , $t_2 = T + \tau/2 + t_1$ and $t_3 = 2T + 3\tau/2 + t_1$ (with $\tau = \pi/\Omega$), respectively, the probability amplitude of excited state is given by

$$c_{2,\mathbf{p}+\hbar\mathbf{k}}(t_3 + \tau/2) = -\frac{i}{2}e^{-i\delta\tau/2}e^{i[\delta t_2 + \phi(t_2)]}(1 - e^{i\delta\tau/2}e^{-i\Delta\phi}) \quad (3.6)$$

where

$$\Delta\phi = \phi(t_1) - 2\phi(t_2) + \phi(t_3). \quad (3.7)$$

In the above equation, $\phi(t_i)$ is the relative phase of the light applying at time t_i , and δ is the two-photon detuning to the transition $|1\rangle \rightarrow |2\rangle$. The atomic population in excited state is given by

$$|c_{2,\mathbf{p}+\hbar\mathbf{k}}(2T + 2\tau)|^2 = \frac{1}{2}[1 - \cos(\Delta\phi - \delta\tau/2)]. \quad (3.8)$$

For the $\pi/2 - \pi - \pi/2$ pulses sequence, if the two-photon detuning δ keeps constant and the phase of Raman pulses is not changed ($\Delta\phi = 0$), a $\pi/2 - \pi - \pi/2$ pulses will bring atoms to ground state $|1\rangle$. According to Equation 3.8, the population of the atom can be tuned by changing the relative phase. If we scan the relative phase of the Raman pulses, the population of atoms vary and the interference fringes can be observed. In our experiment, the phase of

Figure 3.4: A typical Mach–Zehnder-type atom interference fringe by scanning the relative phase of Raman pulses [52]. The interference fringe is obtained with atoms inside the Hollow-core fiber. OD is the optical depth which is defined in section 5.2.

the Raman beam pulses are controlled by an electro-optic modulator (EOM). The Figure 3.4 shows a typical atomic interference fringe as a function of the relative phase through an EOM [52].

3.3 The path integral approach

The phase shifts discussed in this thesis can be sufficiently determined by Equation 3.7. However, this equation is not complete to determine all the phase shifts. From the first principles, a fully general derivation is given by a full quantum mechanical treatment [72]. The total phase shifts can be decomposed into the three terms

$$\Delta\phi_{total} = \Delta\phi_{prop} + \Delta\phi_{sep} + \Delta\phi_{las}. \quad (3.9)$$

The trajectories of the two arms of atoms through the interferometer determine these three components, as determined by the free evolution of the atoms and optical Raman pulses with Lagrangian

$$L = \frac{1}{2}m(\dot{\mathbf{r}} + \boldsymbol{\Omega} \times (\mathbf{r} + \mathbf{R}_e))^2 - m\phi(\mathbf{r} + \mathbf{R}_e) \pm \frac{1}{2}\hbar\alpha\mathbf{B}(\mathbf{r})^2 \quad (3.10)$$

where m is the mass of the atom, $\mathbf{r}$ is the atom's position, $\mathbf{R}_e$ is the radius of Earth, and $\boldsymbol{\Omega}$ is the rotation rate of Earth. $\phi(\mathbf{r} + \mathbf{R}_e) = -(\mathbf{g} \cdot \mathbf{r} + \frac{1}{2!}(T_{ij})r_i r_j + \dots)$ is the gravitational potential, where $\mathbf{g}$ is the gravity and $T_{ij} \equiv \partial_j g_i$ is the gradient tensor of the gravity. $\mathbf{B}$ is the magnetic field that affects the phase shift only through the second-order Zeeman shift, and α is the clock transition Zeeman shift since the atoms are in the Zeeman states $|m_F = 0\rangle$ through the interferometer.

For the $\pi/2 - \pi - \pi/2$ -pulses Mach-Zehnder interferometer, there are five semi-classical interaction vertices (A, B, C, D, and E) through the whole the trajectories (see Fig. 3.5). In the following part, the three components of the total phase shifts will be calculated with respect to these semi-classical interaction vertices.

3.3.1 Propagation phase

The propagation phase shift $\Delta\phi_{prop}$ is from the phase shift picked up by the atom due to the free evolution of atoms wave-packets between the optical Raman pulses. $\Delta\phi_{prop}$ arises from the path integral along the two classical trajectories (one is AB and BD, the other is AC and CE), and depends on both the total energy of atoms and their spatial propagation, including the external energy and internal energy due to different spin states. In the $\pi/2 - \pi - \pi/2$ -pulses Mach-Zehnder interferometer, the pulse sequences is time-symmetric. The two parts of atoms in each of the spin states will travel along the two paths for the same amount of time, and the phase shifts contributing from the internal energy of atoms can be ignored.

In the $\pi/2 - \pi - \pi/2$ -pulses Mach-Zehnder interferometers with the presence of the gravitational acceleration, the propagation phase shift $\Delta\phi_{prop}$ is given by the integral

$$\Delta\phi_{prop} = \frac{1}{\hbar} \oint L[\mathbf{r}(t), \dot{\mathbf{r}}(t)] dt \quad (3.11)$$

Figure 3.5: Space-time diagram for a Mach-Zehnder interferometer in the presence of gravity. The five semi-classical interaction vertices are labeled by ABCDE.

calculated along the two classical paths of the atoms. In the presence of gravitational acceleration with gravity gradients or rotations, the two trajectories will not form a closed loop. For gravity gradients or rotations, the perturbative approach is a useful method, in which $\Delta\phi_{prop}$ is still evaluated along loop formed by the unperturbed trajectories.

3.3.2 Separation phase

In the presence of a constant gravitational acceleration with no gravity gradients and rotations, the two arms of wavepackets would overlap perfectly to form a closed loop at the final beamsplitter. In general, however, the atom interferometer doesn't necessarily form a closed loop. The separation phase $\Delta\phi_{sep}$ is from the imperfect overlap of the two interferometer arms at the end of the interferometer sequence. The two paths may need to originate from different points in space in order to overlap and interfere at the end of the interferometer sequence.

If we assume the initial atomic wave-function is modeled as a plane wave, and the momentum of the atom is $\mathbf{p}$, the phase shift contribution from the spatial separation $\Delta\mathbf{r}$ in the interferometer trajectories is given by

$$\Delta\phi_{sep} = \frac{\mathbf{p} \cdot \Delta\mathbf{r}}{\hbar} = \frac{2\pi\Delta x}{\lambda_{dB}} \quad (3.12)$$

where Δx is the difference of the path length of the two arms of atoms, λ_{dB} is the Broglie wavelength of the atom. It is similar to the phase shift in the optical interferometer.

3.3.3 Laser interaction phase

The laser interaction phase $\Delta\phi_{laser}$ is from the interactions between atoms and light in the two-photon stimulated Raman transitions. When the optical field drives the two-photon transition bringing atoms from one state to another state, the phase of the optical field will be transferred to atoms [73].

In a $\pi/2 - \pi - \pi/2$ pulse sequence, we assume the two-photon detuning $|\delta| \ll \Omega_{eff}$, and the effective Rabi oscillation frequency Ω_{eff} is the same for

all three optical Raman pulses. The transition probability amplitudes of time evolution of the ground and excited state for the case of a $\pi/2$ pulse of length $\tau/2$ is given by

$$c_2(t_0 + \tau/2) = e^{-i\delta\tau/4}[c_2(t_0) - ic_1(t_0)e^{-i(\delta t_0 + \phi)}]/\sqrt{2} \quad (3.13)$$

$$c_1(t_0 + \tau/2) = e^{i\delta\tau/4}[-ic_2(t_0)e^{i(\delta t_0 + \phi)} - c_1(t_0)]/\sqrt{2} \quad (3.14)$$

and a π pulse of length τ is given by

$$c_2(t_0 + \tau) = -ic_1(t_0)e^{-i\delta\tau/2}e^{-i(\delta t_0 + \phi)} \quad (3.15)$$

$$c_1(t_0 + \tau) = -ic_2(t_0)e^{i\delta\tau/2}e^{i(\delta t_0 + \phi)} \quad (3.16)$$

where ϕ is the local phase of the optical field driving the two-photon stimulated Raman transitions.

After a $\pi/2$ - pulse at time $t = t_1$, a π - pulse at time $t_2 = T + \tau/2 + t_1$ and a second $\pi/2$ - pulse at time $t_3 = 2T + 3\tau/2 + t_1$, respectively, the probability amplitude of excited state is given by

$$c_{2,\mathbf{p}+\hbar\mathbf{k}}(t_3 + \tau/2) = -\frac{i}{2}e^{-i\delta\tau/2}e^{i[\delta t_2 + \phi(t_2)]}(1 - e^{i\delta\tau/2}e^{-i\Delta\phi}) \quad (3.17)$$

where

$$\Delta\phi = \phi(t_1) - 2\phi(t_2) + \phi(t_3) \quad (3.18)$$

and

$$\Delta\phi_{laser} = \Delta\phi = \phi(t_1) - 2\phi(t_2) + \phi(t_3) \quad (3.19)$$

is the total differential phase shift due to the interactions between atoms and the three two-photon Raman pulses in a $\pi/2 - \pi - \pi/2$ - pulses Mach-Zehnder interferometer.

3.4 Precision Gravity Measurements

The inertial forces can change the matter-waves' relative phase with respect to the phase of the two-photon Raman optical fields. The phase shift due to the inertial force can be converted to a change in the population of atoms in excited state. Compared to the conventional sensors, the inertial sensors based on atom interferometers, which are isolated from external noises, have a distinguishing feature that a free-falling atoms are used as a reference for laser phases with a sensitivity of laser's wavelength.

We consider the scenario that an atom is driven by a $\pi/2 - \pi - \pi/2$ - pulses Mach-Zehnder interferometer using vertically aligned two counter-propagating laser beams in the presence of an uniform gravitational acceleration g . For constant acceleration and no rotations, $\Delta\phi_{prop} = \Delta\phi_{sep} = 0$, and the differential phase is from the laser interaction phase $\Delta\phi_{laser}$ only.

The phase shift due to the laser interaction given in Eq. 3.19 depends on the phases of the Raman pulses at the time t_1 , $t_2 = T + \tau/2 + t_1$ and $t_3 = 2T + 3\tau/2 + t_1$ respectively. The local laser phase that atoms experience at time t_i is ϕ_i^0 . The laser interaction phase is given by

$$\begin{aligned}\Delta\phi_{laser} &= \phi_1 - 2\phi_2 + \phi_3 \\ &= (\phi_1^0 + k_{eff} \cdot z_1) - 2(\phi_2^0 + k_{eff} \cdot z_2) + (\phi_3^0 + k_{eff} \cdot z_3) \\ &= k_{eff} \cdot (z_1 - 2z_2 + z_3) + (\phi_1^0 - 2\phi_2^0 + \phi_3^0)\end{aligned}\tag{3.20}$$

where the position of the atom is given by

$$z_i = z_0 + v_0 t_i + \frac{1}{2} g t_i^2\tag{3.21}$$

with z_0 is the initial position and v_0 is the initial velocity of atoms. The total phase shift is

$$\begin{aligned}
\Delta\phi &= \Delta\phi_{prop} + \Delta\phi_{sep} + \Delta\phi_{laser} \\
&= \mathbf{k}_{eff} \cdot \mathbf{g}T^2 + (\phi_1^0 - 2\phi_2^0 + \phi_3^0).
\end{aligned} \tag{3.22}$$

Hence the overall phase shift due to the gravity g is given by

$$\Delta\phi_g = -\mathbf{k}_{eff} \cdot \mathbf{g}T^2. \tag{3.23}$$

We can find that this phase shift increases with the momentum transferred N during the beamsplitters and mirrors. In the case that the momentum transferred to atoms in beamsplitters and mirrors is $k_{eff} \approx 2Nk > 2k$, the overall phase shift is

$$\Delta\phi_g = -\mathbf{k}_{eff} \cdot \mathbf{g}T^2 \approx 2N\mathbf{k} \cdot \mathbf{g}T^2. \tag{3.24}$$

By measuring the phase shift $\Delta\phi_g$, and wavevector $\mathbf{k}$ precisely, we can determine the value g with high accuracy from Eq. 3.23.

The phase resolution of the atomic inertial sensors is limited by the atom shot noise. For a sensor with a flux of atoms n , the phase resolution is limited to $\delta\phi \sim 1/\sqrt{n}$. The acceleration sensitivity of the atomic sensor is

$$\delta g = \left(\frac{\delta\phi}{\Delta\phi_g} \right) g = \frac{1/\sqrt{n}}{k_{eff}T^2}. \tag{3.25}$$

The shot-noise-limited sensitivity can be enhanced by increasing the momentum transferred in the beamsplitters and increasing the interferometer time T .

Chapter 4

Experimental apparatus

4.1 Fiber selection

The most conventional fibers are made of silica glass, first developed in the 1960s. The light is confined in the central core by total internal reflection due to higher refractive index of the material than the cladding layer. Recently, a new paradigm in fiber optics has emerged with the ability of optical guidance in air-core by other guidance mechanisms. The air-guidance fibers show the possibility of loading gas particles into the hollow-core fibers. This provides a novel regime to both confine light and gases over long distance. Such regime could create strong enhancement in light-gas interactions and could be used in wide applications ranging from quantum optics to quantum sensing.

The hollow-core photonic crystal fiber (HC-PCF) was first demonstrated experimentally in 1999 [19]. There are two types guiding mechanism of HC-PCF that have developed. The first one is based on the photonic band gap (PBG) in the cladding structure and the second one is related to inhibited coupling between the guiding core modes and the cladding modes.

In the photonic band gap HC-PCF, the cladding structure of the fiber is typically a periodic transverse arrangement of air-holes in a silica matrix and this structure exhibits a PBG. Due to the PBG, light cannot propagate in the cladding structure. If the propagating modes of the core couple with the PBG of the cladding, theses modes are trapped in the core and guided with no

leakage. Using this principle, PBG guiding HC-PCF with an attenuation loss as low as 1.2 dB/km was demonstrated in 2005 [20]. However, the PBG guiding HC-PCF has several limits, such as poor optical power handling, higher group velocity dispersion, and narrow transmission bandwidth.

The second type of HC-PCF is Kagome lattice HC-PCF, which could provide broader bandwidth guidance. However, this HC-PCF has higher transmission loss-levels compared to the PBG guiding HC-PCF. This type of HC-PCF is based on the mechanism of inhibited coupling (IC). It was first reported in 2002 [74], but its guidance mechanism was elucidated in 2007 [75]. Unlike the PBG guiding HC-PCF, the cladding in the IC guiding HC-PCF doesn't exhibit any PBG but relies on a strong transverse-field mismatch between the continuum of core modes and cladding modes. The IC guiding HC-PCF exhibits loss figures of only a few tens of dB/km loss level with a record loss level of 17 dB/km at 1064 nm demonstrated in 2013 [76].

In our experiment, we use a hollow-core photonic crystal fiber with the model number *PMC-C-TiSa-Er-7C* from GLOphotonics. The fiber has a low transmission loss at 800 nm and 1600 nm wavelength. At 800 nm, the low loss transmission window covers rubidium D_1 and D_2 transitions at 795 nm and 780 nm, respectively, see Fig. 4.1). The attenuation at the wavelength of 795 nm and 780 nm is lower than 80 dB/km which is negligible for the 4 cm fiber length used in the experiments in this thesis.

The inner core diameter of the fiber is $63 \pm 1 \mu\text{m}$. The mode field diameter is $44 \pm 1 \mu\text{m}$, and it is nearly a Gaussian mode in the far field. The distance between the wall and the $1/e$ distance of the fundamental mode is about $10 \mu\text{m}$. Therefore, the interaction between the fiber wall and atoms inside the fiber can be ignored. Before we assemble the hollow-core fibers into the ultra-high vacuum chamber, we check the mode profile of the fibers to ensure a Gaussian-like mode by using a CCD camera. For the purpose of atom interferometers experiments in the fiber, the choice of the fiber core size is really a tradeoff between atom numbers and optical power required. Small core size fiber, such as the $5.5 \mu\text{m}$ core diameter from Thorlabs of model number HC-800B, would results in small number of atoms loaded into and fiber and stronger surface

Figure 4.1: (a) Data of attenuation and dispersion spectrum of the $PMC - C - TiSa_{Er} - 7C$ fiber used in our experiment [78]. (b) Typical output near field image at 800 nm from the vendor [77].

interaction. On the other hand, larger core size will require more optical power for trapping and interferometer pulse, and higher order modes could also affect the performance. The more details about the fiber selection are discussed in Appendix B.

4.2 Vacuum system

Our vacuum system mainly consists of an 8.0'' spherical octagon chamber from Kimball Physics Inc., a fiber holder for mounting the hollow-core fibers, a Rubidium ampoule glass-sealed in Argon, an ion pressure gauge, and an ion pump (see Fig. 4.2). The viewports are anti-reflection coated for high transmission at near-infrared wavelength light used in our experiment. Viewports, vacuum flanges, and fiber holder were washed by ultrasonic bath before assembly, for details, see Appendix C. After the main vacuum chamber and the fiber holder were assembled, we use an external turbo pumping system to pump the chamber for a day and the pressure of the vacuum chamber went down to 1×10^{-5} torr. We then heated the entire chamber at 120 °C for two weeks while keeping the turbo pump on only. The ion pump was turned on and the baking lasted for one more week. The pressure of the chamber can achieve to 1×10^{-9}

Figure 4.2: The vacuum chamber used in our experiment. The fiber holder designed by ourselves is used for mounting hollow-core fiber.

torr. We then turned off the baking and closed the all metal gate valve on the chamber before the turbo pump was switched off. The ion pump alone can maintain the pressure at the level of 10^{-9} torr before releasing rubidium atoms into the vacuum chamber. After the rubidium ampoule were broken, the pressure inside the chamber went up to about 5×10^{-9} torr.

The fiber holder was designed to hold three fibers and is made of aluminum (see Fig. 4.3). In the experiment, we found that the aluminum would block the microwave signal. For the next version, we will design the fiber holder using ceramic. In order to ensure there is enough thermal contact between the fibers and the holder for baking, we put a layer of high conductivity thermal paste (Silver-Filled paste from Kurt J. Lesker Company) between the fibers and the holder. At this stage, we were not sure about the pressure inside the fiber since the fiber with a small core is a long differential pumping tube and the outgassing rate on the fiber inner surface is not known. Although

Figure 4.3: Design of fiber holder used for mounting the hollow-core photonic crystal fibers with vacuum compatible glue. We also put the thermal paste between the fibers and holder to increase the heat conductance for baking.

there were some studies regarding the pressure inside a hollow-core fiber [28], the outgassing rate on the inner core surface actually depends on the manufacturing and handling process. Cooling the fiber to liquid-nitrogen or even liquid-helium temperatures [78] would resolve the problem, but it would make our system more complexing. As a first step, we baked our chamber with care for a little more than a month.

For the first iteration of the experiment, we placed a rubidium dispenser in the vacuum chamber for the rubidium source. The dispenser was controlled by connecting pins on an electrical feedthrough. We ran few amperes of current through the dispenser for 5 minutes and then we turned down the current to a level that will keep the pressure of the chamber constant throughout the day.

The pressure of the chamber can stay constant just by slightly adjusting the current to the dispenser few times during a day. After the rubidium atoms were introduced into the chamber, the pressure was about 5×10^{-9} torr. We ran out of the dispenser after two years of operation and replaced the rubidium dispenser by a rubidium ampoule.

4.3 Laser system

Figure 4.4: ^{85}Rb D_2 line (780 nm) hyperfine structure and transitions for cooling, repumping, probe, push and Raman beams.

In this section, we introduce all the laser systems used in the MOT, dipole trap, and detection. Figure 4.4 shows the laser beams and energy transition used in our experiment.

4.3.1 Cooling lasers

Figure 4.5: Optical setup of the laser beams used in the experiment in an optical SMART table.

The master laser beam is locked directly to the transition of $F = 3$ to $F' = 2/4$ crossover (offset-locked -92 MHz to $F = 3 - F' = 4$ transition) using saturated absorption spectroscopy. The cooling and imaging beams are derived from the master beam. The master beam is shifted up 80 MHz by one double-pass acousto-optic modulator (AOM) and one single pass AOM setup ($+2 \times 80$ MHz $- 80$ MHz). The frequency of the cooling beam is tuned by the double-pass AOM to 12 MHz below the $F = 3$ to $F' = 4$ transition. The single pass AOM is used for switching the cooling beam. The cooling beam is also used for imaging (correspond to $F = 3$ to $F' = 4$ transition) the MOT by shift the AOM frequency ($+2 \times 86$ MHz $- 80$ MHz).

The repumper beam is directly locked to the transition of $F = 2$ to $F' = 1/2$ crossover using another saturated absorption spectroscopy. It is then shifted 80 MHz up at be the transition of $F = 2$ to $F' = 3$ by a single pass AOM. The repumping beam is used to bring atoms from the $F = 2$ state back to the $F = 3$ state for continuation of cooling process.

4.3.2 State preparation and detection Lasers

Figure 4.6: The optical setup for the MOT and detection. (CYL: cylindrical lens pair, HWP: half-wave plate, QWP: quarter-wave plate, PBS: polarizing beam splitter, the colored boxes indicate the fiber coupler.)

For preparing atoms in the $F = 2$ state, we use the depump beam to bring atoms from ground state $F = 3$ to ground state $F = 2$. The beam is locked to the transition $F = 3$ to $F' = 2/4$ crossover; then it is shifted by an AOM to the resonant transition $F = 3$ to $F' = 4$.

A push beam is used to bring atoms from the $F = 3$ state to $F = 2$ state (dark state). The push beam is directly locked to the transition of $F = 3$ to $F' = 2/4$ crossover; then is shifted 152 MHz to the resonant transition $F = 3$ to $F' = 4$ by a double-pass AOM and a single pass AOM. The push and repumping beams can work together to perform the $F = 3$ to $F' = 4$ cycling transition, which is used to push away atoms outside fiber before detection.

The probe beam for detecting atoms is generated by shifting the master beam 92 MHz up to the transition $F = 3$ to $F' = 4$ by one double-pass AOM and one single pass AOM setup ($+2 \times 86 \text{ MHz} - 80 \text{ MHz}$). The detuning of the probe beam can be shifted by the double pass AOM for absorption measurement.

All lasers used for MOT and detection are based on external cavity laser diodes [79]. The optical setup for the MOT and detection beams can be seen in the Fig. 4.6. The cooling and repumping beam are coupled into the fibers for mode shaping. The probe beam is picked off from the zero-order of the cooling beam by using a D-shaped mirror. By shifting the master beam 160 MHz up by a double-pass AOM ($+2 \times 80 \text{ MHz}$), the laser from the diode is used for the injection beam (45 mW) of a tapered amplifier (TA). After the TA, the beam is shifted 80 MHz down by a single pass AOM. Hence, the cooling beam is 12 MHz below the $F = 3$ to $F' = 4$ transition. The total power of the cooling beam is about 60 mW. It is divided into three beams for MOT building. They are then coupled into three fibers with 20 mW of power in each port. The repumping power used for MOT is about 7 mW.

4.3.3 Dipole laser

A red-detuned (red-detuned from the D_1 line transition-795 nm) dipole trap is used for loading atoms into the hollow-core fiber and confining the atoms inside the hollow-core fiber from the collision to the fiber inner wall. The dipole beam is a single-mode beam guided by the hollow-core fiber.

The dipole laser is from a Titanium-sapphire Laser bought from M Squared. Figure 4.7(a) shows the working principle of the Titanium-sapphire Laser. The Titanium-Sapphire refers to the lasing medium, which is a crystal of sapphire

doped with titanium ions. Another laser with a wavelength of 532 nm is used to pump the crystal of sapphire to generate light with wavelengths near 800 nm. In our experiment, the wavelength of the dipole beam can be tuned at the range of 808 nm-875 nm by changing the birefringent filter. Laser beam from the Ti-Sapphire Laser is coupled to a single mode fiber in order to decouple it from the rest of the optical alignment. The total power of the beam and the coupling efficiency as a function of the pumping power are shown in Fig. 4.7(b). The coupling efficiency would slightly decrease at high optical power.

Figure 4.7: (a) The total power of the laser beam as a function of pumping power of the Ti-Sapphire laser. (b) The coupling efficiency of the beam to a single mode fiber.

4.3.4 Raman Lasers

The two lasers with different frequencies for the two-photon Raman transition is generated from a single laser passing through an EOM. The frequency of the laser is locked to ^{87}Rb D_2 line from $F = 2$ to $F' = 2/3$ crossover by saturated absorption spectroscopy, which is 1.32 GHz red-detuned from the resonant transition $F = 3$ to $F' = 3$ (see Fig. 4.4). It is then shifted 200 MHz up to 1.34 GHz detuned to the transition $F = 3$ to $F' = 3$ by an AOM for switching. The EOM is frequency modulated at 3.04 GHz, the transition of ^{85}Rb ground states.

The electric field of laser light before the EOM is

$$E_{\text{before}} = E_0 e^{i\omega_L t}. \quad (4.1)$$

After the EOM with modulation frequency $\Omega_M = 3.04$ GHz and the modulation amplitude of β , the electric field is

$$E_{\text{after}} = E_0 e^{i\omega_L t + i\beta \sin(\Omega_M t)} = E_0 e^{i\omega_L t} \sum_{k=-\infty}^{\infty} J_k(\beta) e^{ik\Omega_M t} \quad (4.2)$$

where $J_k(\beta)$ is the Bessel function. For $\beta \ll 1$, only zero and first orders are considered

$$E_{\text{after}} = E_0 \left(e^{i\omega_L t} + \frac{\beta}{2} e^{i(\omega_L + \Omega_M)t} - \frac{\beta}{2} e^{i(\omega_L - \Omega_M)t} \right). \quad (4.3)$$

The three components are carrier, +1 sideband and -1 sideband respectively. Hence, the frequencies for the two photon Raman transitions are formed by the carrier and first sideband of the laser after the EOM. By adjusting the power ratio between the carrier and sideband, the ac Stark shift introduced by the Raman beams can be minimized [80]. We note that residual ac stark shift from the higher order sidebands beams can be removed by using two independent lasers phase-locked at 3.04 GHz.

4.3.5 Optical pumping laser

Figure 4.8: (a) Energy level for optical pumping laser beams. (b) Diagram of the selection rules of the transitions using the π -polarized beam. To increase the S/N ratio, we introduced the optical pumping technique to our experiment. The transition $m_F = 0 \rightarrow m_{F'} = 0$ is a forbidden transition, resulting in $m_F = 0$ to be the dark state where atoms are pumped to.

In the presence of an external magnetic field, the energy of ground state $F = 2$ is split to five Zeeman substates. Atoms in the ground state $F = 2$ will distribute among all five Zeeman substates. However, only atoms in the clock state ($m_F = 0$) will be driven by the two-photon Raman pulses in the atom interferometer, resulting in only 20% of atom number used in the experiment.

In order to increase the atom numbers for the experiment, the optical pumping process is employed. We use a π -polarized laser beam (the quantization axis is aligned along the direction of the electric field) resonant on $F = 2 \rightarrow F' = 2$ transition. In this case, the transition $m_F = 0 \rightarrow m_{F'} = 0$ is forbidden. Once atoms are pumped to the $m_F = 0$ state, they will not be driven by the laser beam. Atoms also have finite probability to be pumped into $F = 3$ state. We apply a repumping beam to recover those atoms back to $F = 2$ state. In principle, this will increase the signal to noise ratio of the atom interferometer experiment by a factor of 2.24. The working principle is shown in Fig. 4.8.

Figure 4.9: (a) The main vacuum chamber in the experiment. (b) An image of cold atomic cloud collected in the MOT. The red spot in the center of the picture is the atomic cloud.

4.4 Laser cooled atomic sources

A three-dimensional MOT collects rubidium atoms from background atomic vapor in vacuum. The setup of the MOT consists of three retro-reflected laser beams (the beams are shown in the optical setup in previous section) and a three-dimensional quadrupole magnetic field produced by anti-Helmholtz coils are described in Fig. 4.9. There are also three pairs of auxiliary coils in our MOT setup to compensate the stray magnetic fields. They are used to shift the MOT position which are useful for loading atoms. A CCD camera is mounted for imaging atoms. It is used to monitor the position of MOT as well as measure the MOT temperature through TOF. A picture of an ensemble of atoms from the CCD camera is shown in Fig. 4.10 (b). The diameter of the cloud is around 2 mm. The cloud is shifted exactly on top of the hollow-core fiber by the three pairs of auxiliary coils.

Eq. 2.35 is used to determine the temperature of the cloud through the measurement of the ballistic expansion of the cloud over time. Density distribution of the atom cloud as a function of the position on the z -axis at $x = 0$,

Figure 4.10: Distribution of the density as a function of position with temperature T at different releasing time.

$y = 0$ at different times is shown in Fig. 4.10.

We fit the curves by a Gaussian function, and extract the $1/e$ radius as σ_t . We plot the $\sigma_t^2 - t^2$ graph and fit it according to the expression

$$\sigma_t = \sqrt{\sigma_0^2 + \sigma_v^2 t^2}. \quad (4.4)$$

The Gaussian radius σ_v of the velocity distribution is associated with the temperature T of the cloud by the formula

$$T = \frac{M}{k_B} \sigma_v^2. \quad (4.5)$$

According to the linear fitting in Fig. 4.11, the temperature of the atoms is $\sim 10 \mu\text{K}$.

Figure 4.11: Standard deviation of the Gaussian fit as a function of t^2 , where t is the releasing time of the atomic cloud.

4.5 Control system

The computer-based control system is responsible for timing sequence initialization, power controlling, data acquisition, diagnostics, signal processing, and frequency synthesis. The laser system is also controlled by this system, including the temperature and current control, frequency stabilization and parameter of the laser beams (power and frequency). A schematic of the control system for the laser beams is shown in Fig. 4.12.

All the instruments in our experiment are controlled with output from National Instruments boards PXIe-6356. It is a multifunction data acquisition device, and it offers analog Input/Output, digital Input/Output(I/O), and

Figure 4.12: A schematic of the control system for all the laser beams used in the experiment.

digital TTL pulses. The digital pulses are used for switching the EOMs and AOMs. The analog outputs are used to control the optical power of the laser beams through controlling the RF power of the AOMs. They are also used to control the magnetic field coils.

Few programmable direct digital synthesizers (DDSs) are used to generate variable RF frequencies required for the EOMs and AOMs. The model number of the DDSs used is AD9959. The detuning frequency of the cooling and repump beams are controlled by the DDS boards. They are also used to adjust the frequency of the microwave signal, and the velocity of the moving optical lattice by jumping or ramping the frequencies of the dipole beams. The relative phase of two-photon Raman beams is controlled by the phase modulation of the DDSs.

4.6 Microwave frequency generation

To drive transitions between the two hyperfine ground states of the atom, a microwave signal of frequency 3.04 GHz is used. The EOM is also driven by the microwave signal to generate the two-photon Raman beams and the relative phase of the Raman beams is controlled by the microwave source too. Our low phase noise microwave source follows the design in reference [81] as shown in the Fig. 4.13.

A 10 MHz atomic lock (FS725) with low frequency noise from SRS provides the clock reference to the frequency generator. The frequency generator provides the LO input for the single sideband modulator with model number HMC496LP3. A direct digital synthesis (DDS) board provides the input signals of the I and Q of the modulator, which are used to provide the coherent phase control on the frequency. The input signals can control the phase and amplitude of the RF signal for the microwave. Then the microwave signal is sent to the high power amplifier by a low loss cable CA400. The microwave signal after the amplifier is sent to a microwave antenna.

Figure 4.13: A schematic of the design for the low phase noise microwave source.

4.7 Detection devices

We use Andor camera of model number DU-897U-CS0-EXF for taking the fluorescence images of atoms inside the fiber. We use it to monitor the loading process of atoms as well. By using this camera, we easily can achieve high quality and faster time-lapse imaging.

A CCD camera of model number acA640-90um-Basler-ace is used to take the fluorescence imaging of the atomic cloud collected in the MOT. The time-of-flight images are also implemented by the CCD camera. The acA640-90um-Basler-ace is a USB 3.0 camera with the sensor with mode of Sony ICX424 CCD. It delivers ninety frames per second at VGA resolution. The basic information of the CCD camera is shown in Fig. 4.14. The resolution of the camera is enough for the fluorescence imaging of the atomic cloud.

Basic Information	
Vendor	Basler
Model	acA640-90um
Type of data presented	Typical
Number of samples	100
Sensor	Sony ICX 424
Sensor type	CCD
Sensor diagonal	Diagonal 6 mm , Optical Size 1/3"
Indication of lens category to be used	C-Mount
Resolution	659 x 490 pixel
Pixel width	7.40 μm
Pixel height	7.40 μm
Readout type	Progressive scan
Transfer type	Interline transfer
Shutter type	-
Overlap capabilities	Overlapping
Maximum readout rate	90 frames/second
General conventions	-
Interface type	USB 3

Figure 4.14: Basic information of the CCD camera acA640-90um-Basler-ace.

For the detection of the probe beam through the hollow-core fiber, we used a high sensitive APD module. The model number of the APD module is C10508-01 from Hamamatsu. The APD module consists of an APD sensor and a converter (convert current signal to voltage signal). The micro-controller for compensating temperature and a voltage power supply circuit are also provided. The specifications of the APD is shown in Fig. 4.15.

Type	High-stability type
Effective photosensitive area	$\phi 1.0$ mm
Temperature stability of gain (typ.)	- %
Built-in sensor	S12023-10A
Photoelectric sensitivity (typ.)	1.25×10^7 V/W
Cooling	Non-cooled
Low Band Cut-off Frequency (typ.)	DC MHz
High Band Cut-off Frequency (typ.)	10 MHz
Supply voltage (typ.)	± 5 V
Measurement condition	Photoelectric sensitivity: $M=250$, $\lambda=800$ nm, Temperature stability of gain: $0\sim 40$ °C

Figure 4.15: Specifications of the APD with C10508-01 bought from Hamamatsu. The photoelectric sensitivity is 1.25×10^7 V/W.

Chapter 5

Cold atoms inside a hollow-core optical fiber

5.1 Fiber preparation

We sliced a few pieces of 4-cm-long fibers and strip their coating manually by a sharp razor blade. The image of the fiber end under a microscope after slice is shown in Fig. 5.1(a). A beam profiler is used to check the far field beam profile as shown in Fig. 5.1.

To couple the laser beams into the fiber, we use a lens with focal length ~ 18 mm to collimate the beam from a single mode fiber. A lens with focal length ~ 150 mm is then applied to focus the beam into the fiber. The coupling efficiency to the hollow-core fiber is about 85%. The beam through the hollow-core fiber is then coupled to another single mode fiber. The coupling efficiency to the single mode fiber is 70%. We put three fibers in the chamber for the experiment with around 80% coupling efficiency for each fiber. The beam profiles of each fiber are checked to be nearly Gaussian.

We use the knife edge measurement to characterize the properties of the fiber mode. The power distribution of the beam is measured at different z position along the fiber axis. The data is shown in Fig. 5.2. The beam waist w is extracted from the data of power distribution. Then plot the $w(z) - z$ curve and fit it using the equation

Figure 5.1: (a) Image of the hollow-core fiber we used under a microscope. (b) Typical near field intensity profile of the beam coming out of the fiber.

Figure 5.2: (a) Power distribution of the beam along the transverse direction from the hollow-core fiber. (b) The beam waist at different positions in z -axis along the fiber.

$$w(z) = \sqrt{w_0^2 + (\phi z)^2} \quad (5.1)$$

where

$$\phi = \frac{\lambda M^2}{\pi w_0}. \quad (5.2)$$

The beam waist measurement results are

The Rayleigh length $z_0 = (8.29 \pm 0.09)$ mm

The size of the mode inside the fiber $w_0 = (0.0233 \pm 0.0005)$ mm

The beam quality factor $M^2 = 1.95 \pm 0.08$

M^2 value of close to 2 is typical of a fiber tip output.

5.2 Atoms inside the fiber

The temperature of the atoms is measured to using TOF method to ~ 10 μ K. We also measure the number of the atoms by the fluorescence imaging method. The estimated number of atoms is about 5.6×10^8 with the density of $8.6 \times 10^9/\text{cm}$.

Figure 5.3: The schematics of the experimental setup for loading and probing of atoms.

To load and confine atoms in the fiber, a dipole beam with red detuning frequency is coupled into the hollow-core fiber, see Fig. 5.3. The wavelength of the dipole beam can be adjusted from 808 nm to 875 nm and the power is about 500 mW inside the fiber. Due to the deep trapping potential (i.e. strong transverse confinement) given by the guided mode of the dipole beam inside fiber, the atoms are radially confined inside the fiber. Cold atoms used for the loading are a laser-cooled ^{85}Rb atoms collected on top of the fiber. They are then transferred into the dipole trap. The potential gradient created by the diverging beam from the fiber then attracts cold atoms into the fiber core.

Figure 5.4: A 3D schematic of the loading setup and the arrangement of the MOT beams relative to the fiber location.

To detect the atoms inside hollow-core fiber, we measure the transmission of a probe beam with an avalanche photodiode (APD). The probe beam is coupled into the hollow-core fiber together with the dipole beam from the bottom of the fiber output, as shown in Fig. 5.3. After the hollow-core fiber, they are collimated by a lens outside the chamber and then coupled into a single mode fiber. The single mode fiber can spatially filter the probe beam after the hollow-core fiber which can filter out the higher order modes. A notch filter and a PBS are used to separate the probe beam from the dipole beam. Finally, the probe beam is detected by the APD.

We demonstrate an efficient method to transfer MOT to the dipole trap inside the fiber. In this method, the three pairs of auxiliary coils in our MOT setup plays an important role in the transfer process. At the end of the MOT loading phase, the current of z -coil are tuned to shift the MOT closer to the facet of the fiber tip, and the x -coil and y -coil are used to shift the MOT in xy -position to maximize the overlap between the ensemble and the dipole beam area. Cold atoms were collected by the MOT for one second. The sub-Doppler cooling is performed by reducing cooling and repumping power to one-third of MOT power, changing the detuning of cooling beam, and turning off the main magnetic field at $t = 100$ ms. At $t = 101$ ms, we adjust the current in the z -coil to shift the cloud closer to the fiber tip for optical dipole trap to take over. Figure 5.5 are the fluorescence images of the atomic cloud. The figure shows the positions of atomic cloud at different times during the loading process, where MOT was released at 101 ms. After ~ 20 ms of free fall, atoms start to move into the fiber. Most atoms are trapped into the fiber after ~ 30 ms.

Figure 5.5: The positions of atomic cloud at different releasing times which starts at 100 ms. After releasing for ~ 20 ms, atoms start to get loaded into the fiber. The end of the loading process is about ~ 30 ms after release.

Figure 5.6: OD as a function of the probe time with different optical dipole beam wavelength.

Figure 5.7: Transmission of probe beam as a function of detuning frequency of the probe beam.

After loading, the proof of atoms trapped inside fiber should be given. We use absorption of a probe beam as a method to characterize the number of

atoms inside the fiber. The transmission of the probe beam is related to optical depth $OD = -\ln(I_{out}/I_{in})$, where I_{in} is the intensity of probe without atoms and I_{out} is the intensity with atoms. The data of the OD as a function of releasing time is shown in Fig. 5.6. Lifetime of atoms inside the fiber is ~ 20 ms. Optical depth inside fiber has been observed up to ~ 200 . We also obtain the data of transmission as a function of detuning of the probe beam with the OD ~ 100 (see Fig. 5.7).

5.3 Number of atoms inside fiber

To estimate the number of atoms inside the fiber, the probe beam resonant with the $F = 3$ to $F' = 4$ transition is sent into the fiber and the transmission T_r of the probe beam is measured. The number of atoms inside the fiber can be determined by $OD = -\ln(T_r) = n\sigma L$, where L is the length of the hollow-core fiber, n is the density of atoms, and σ is the scattering cross section of atoms.

The optical power of probe beam is 3 nW, and the duration of the probe pulse is 50 μ s. We sent two probe pulses, one for detecting atoms and the other one for reference. The second pulse is sent 30 ms after the first pulse to ensure that atoms diffuse out of the probe beam area. The probe power is chosen to avoid saturating atoms because the saturation power, averaging over the transitions, is about 30 nW. To avoid any transient effects of the pulses, we selected only the middle 40 μ m of the pulses for OD measurement. The probe beam is linearly polarized and 8.3 MHz blue-detuned from $F = 3$ to $F' = 4$. It is then detected by an APD (C10508-01).

Assuming the density distribution of atoms is Gaussian and integrating over the transverse profile of the fiber, we estimated the number of atoms loaded into the fiber by integrating

$$\begin{aligned}
 OD &= ((2/\pi W^2))L \int_0^R n_0 \exp(-r^2/r_0^2) \sigma \exp(-2r^2/W^2) 2\pi r dr \\
 &= (2\sigma n_0 L r_0^2 / (W^2 + 2r_0^2)) (1 - \exp(-R^2(W^2 + 2r_0^2)/(W r_0)^2)) \quad (5.3)
 \end{aligned}$$

where $\sigma = \sigma_0/(1 + 4(\delta_p/\Gamma)^2)$ is the scattering cross section of atoms at the center of the dipole trap, σ_0 is the resonant cross section, δ_p is the detuning of the probe beam, $\Gamma = 6.1$ MHz is the decay rate of the excited state, r_0 is the 1/e radius of the initial atomic cloud size, r is the radial coordinate, R is the radius of the fiber core, and n_0 is the peak atomic density. The radius of the atomic cloud can be inferred from the ratio of the cloud temperature T_p and trapping potential U_0 as $r_0^2 = (k_B T_p / (2U_0)) W^2$. For $T_p = 23$ mK, $U_0/k_B = 0.53$ mK, the linearly polarized probe beam resonant on $m = 0$ to $m' = 0$ Zeeman transition, and $\sigma_0 = 0.53 \times 10^{-9}$ cm², arrived at by averaging over all allowed transitions between $F = 3$ to $F' = 4$ Zeeman states, OD = 1 is approximately equivalent to 1.5×10^4 atoms for resonant linearly polarized probe beam.

Chapter 6

Atom interferometry inside a hollow-core photonic crystal fiber

In highly sensitive light-pulse atom interferometers, matter-wave beam splitters and mirrors are formed by optical pulses along the atoms' trajectories. Similar to Ramsey separated oscillatory field method [82], the beam splitter and mirror optical pulses serve as local oscillators to compare the phase difference between two interferometer arms via atom-light interaction. In free space, light-pulse atom interferometers are versatile tools for studying fundamental physics [83–86] and performing inertial sensing [87]. They have leverage over classical sensors because of the universal properties of atoms and long-term stability [88]. Nonetheless, the unavoidable bulky and power-consuming designs make it very difficult to bring atoms close to surfaces or source masses to maximize signals and thus have limited their use in applications requiring small-scale devices or high spatial resolution. For example, to ensure a flat optical wavefront and uniform atom-light interaction, an interferometer beam waist of 2 cm is used in a 10-m-long interferometer, which requires 30 W of optical power at 30-GHz detuning from the resonance [38]. Guiding atoms by light in a micrometer-scale photonic crystal fiber can overcome these obstacles and could enable remote measurement at a targeted location.

Over the past two decades, rapid progress on the design and manufacture of hollow-core photonic crystal fibers provides a nearly single mode and diffraction-free optical field for enhancing atom-light interaction such that the confined optical fields can be used for trapping, guiding, and manipulating atoms. Spectroscopic experiments have shown that thermal or cold atoms can be confined or optically guided in hollow-core photonic crystal fibers for quantum and nonlinear optics applications. However, the requirement of the optical wavefront of the fiber for atom interferometer applications is more stringent, where the coherence of optical field in time and spatial domains are crucial for manipulating quantum states of atoms. Here, we optically trap cold ^{85}Rb atoms in a hollow-core photonic crystal fiber with an $1/e^2$ mode field radius of $22\ \mu\text{m}$ and use the waveguide fields as matter-wave beam splitter and mirror pulses to demonstrate that the coherence of a quantum superposition state of atoms can be preserved and interrogated by the optical guided mode to form an interferometer sensitive to the gravity.

6.1 Temperature of the atoms inside the fiber

The loading process of atoms is presented in chapter 5. Here, I simplify the process as follows. Firstly, cold Rubidium-atoms are prepared on top of the fiber by a three dimensional MOT. Secondly, atoms are released and an external magnetic fields along x , y , z directions are adjusted to bring the atomic ensemble closer to the fiber tip and maximize the number of atoms loaded into the fiber. Finally, when atoms are close to the tip of the fiber, the optical dipole trap and the gravity take over to attract atoms into the fiber.

The experimental loading setup is shown in Fig. 6.1. A 4-cm-long hollow-core fiber is mounted vertically on the fiber holder inside the vacuum chamber. The cold ^{85}Rb atoms are prepared about 0.5 cm above the fiber tip. After releasing atoms from the MOT at $t = 0$, we guide atoms into the fiber by an optical dipole force from the divergent profile of optical field at the facet of the fiber. The frequency of the dipole beam is tuned below the atomic transition frequency, in which atoms are attracted toward higher intensity locations (the center of the fiber). The wavelength of the dipole beam is 808 nm, and the

optical power of the beam is about 280 mW. The calculated trap depth of the beam is 0.53 mK, assuming the dipole beam waist is 22 μm . In addition, we reduced the magnetic field along the vertical axis from 1.90 to 0.61 G to transport atoms closer to the tip of the fiber to improve the loading efficiency. We sent the push beams with 2-ms duration resonant on both the $F = 3$ to $F' = 4$ and $F = 2$ to $F' = 3$ transition to push atoms above the fiber away from the dipole beam area. The push pulse is 5 ms before we probe atoms at $t = t_p$.

Figure 6.2: Measurement of atom temperature T_p in the radial direction at different probe times. The release time in the x axis is the duration where the dipole trap is off or atoms to expand. The error bars indicate the standard error of seven experimental runs. The inset is a plot of temperature at different probe times.

When atoms are inside the fiber, the optical dipole trap is switched off, allowing the atomic cloud to expand freely. After a release time, we sent the

probe beams for OD measurement. In Fig. 6.2, we measure OD versus the time after releasing atoms from the optical dipole beam and obtain the temperature T_p of the atomic cloud in the radial direction at different probe times from the fitted curves. We considered the ballistic expansion of the atomic cloud and thus replaced r_0^2 by $r^2(t) = r_0^2 + v_p^2 t^2$ in Eq. 6.1. To reduce the parameters for fitting, we further used $v_p^2 = 2k_B T_p / m$ and $r_0^2 = W^2 k_B T_p / 2 / U_0$ in Eq. 6.1. With $U_0 = 0.53$ mK, $W = 22$ mm, and $R = 31.5$ mm used in the experiment, we fit the data with an equation

$$\text{OD}(t) = \text{OD}_0 (1 - \exp(-C(1/(AT_p + BT_p t^2) + 1))) \times (1/(AT_p + BT_p t^2 + 1)) \quad (6.1)$$

where OD_0 is the optical depth when the dipole trap is turned off, $A = k_B / U_0 = 0.0019 \text{ mK}^{-1}$, $B = 4k_B / m / W^2 = 0.79 \text{ mK}^{-1} \text{ms}^{-2}$, and $C = 2R^2 / W^2 = 4.1$. The fitting results are shown in Fig. 6.2. The measurement shows an exponential increase of atom temperature over time, which can be explained by heating due to the intensity noise of the dipole beam [89].

6.2 Atom interferometer inside fiber

Our interferometer pulses are generated from a diode laser locked to $F = 2$ to $F' = 3$ of ^{87}Rb , which is $\Delta = 1.34$ GHz red-detuned from the ^{85}Rb $F = 3$ to $F' = 3$ transition. The Raman fields are formed by a fiber EOM modulated around ^{85}Rb ground-state hyperfine splitting frequency. However, 0-order and +1-order of the fields after EOM are used for the two-photon Raman transition. Although 0-order and -1-order also form another pair of two-photon transition with opposite phase. The effective two-photon Rabi frequency at resonance is

$$\Omega_R = \sqrt{\Omega_{R1}^2 + \Omega_{R2}^2 + 2\Omega_{R1}\Omega_{R1} \cos \bar{\varphi}} \quad (6.2)$$

where $\Omega_{R1} = \frac{\Omega_1 \Omega_0}{2\Delta}$ is the two-photon Rabi frequency of the 0-order and the +1-order of the fields, and $\Omega_{R2} = \frac{\Omega_{-1} \Omega_0}{2(\Delta + w_{hfs})}$ is the two-photon Rabi frequency of the 0-order and the -1-order of the fields. The phase term $\bar{\varphi} = \varphi_1 + \varphi_{-1} - 2\varphi_0$,

where φ_i is the phase of 0-order, +1-order, and -1-order of the fields, and $w_{hfs} = 3.036$ GHz is the energy splitting of hyperfine ground states of ^{85}Rb . The EOM produces sidebands with opposite phases ($\bar{\varphi} = \pi$). Thus, the effective two-photon Rabi frequency at resonance is calculated

$$\begin{aligned}
 \Omega_R &= \sqrt{\Omega_{R1}^2 + \Omega_{R2}^2 + 2\Omega_{R1}\Omega_{R2} \cos \pi} \\
 &= \sqrt{\Omega_{R1}^2 + \Omega_{R2}^2 - 2\Omega_{R1}\Omega_{R2}} \\
 &= \frac{w_{hfs}}{(\Delta + w_{hfs})} \frac{\Omega_1 \Omega_0}{2\Delta} \\
 &= \frac{w_{hfs}}{(\Delta + w_{hfs})} \Omega_{R1}.
 \end{aligned} \tag{6.3}$$

We can define the reduction factor by

$$\eta = \frac{w_{hf}}{(\Delta + w_{hfs})}. \tag{6.4}$$

Hence, its the single-photon detuning in our experiment will give a reduction factor of $\eta = w_{hfs}/(1.34 \text{ GHz} + w_{hfs}) \approx 0.69$ on the effective two-photon Rabi frequency [90]. The fields after passing through EOM are split into two optical paths and coupled into the hollow-core fiber from both ends. We further tuned the EOM modulation frequency to match the Doppler shift of atoms. The polarization of the counter-propagating fields is linear and orthogonal to each other. The optical power on the top and bottom Raman beams is 4 and 8 mW, respectively. We shifted the phase of the microwave sent to the EOM during the last interferometer pulse for scanning the interference fringes. The probe beam in the interferometric measurement is resonant on $F = 3$ to $F' = 4$.

To demonstrate an atom interferometer, we run a Mach-Zehnder sequence comprising a beam splitter ($\pi/2$) pulse followed by a mirror (π) pulse and another beam splitter ($\pi/2$) pulse. The timing sequence of the interferometer is shown in Fig. 6.1(b). Atoms are initially prepared in the $F = 2$ state by a depumping beam on $F = 3$ to $F' = 3$ transition for 1 ms. These mirror and beam splitter pulses are formed by a pair of counter-propagating beams

driving Doppler-sensitive two-photon Raman transitions between $F = 2$ and $F = 3$ states. The detuning frequency of the Raman beams is 1.34 GHz red-detuned from the ^{85}Rb $F = 3$ to $F' = 3$ transition. With this detuning, we only use about 10^8 photons for beam splitter pulses. An external magnetic field about 3.6 G is applied along the z direction as a quantization axis to break the degeneracy of Zeeman states such that only atoms projected into $m_F = 0$ Zeeman state interact with interferometer pulses.

After the last pulse, the transition probability of atoms in the $F = 3$ state is

$$(1 - \cos(\Delta\phi))/2 \quad (6.5)$$

where

$$\Delta\phi = \phi_1(t_0) - 2\phi_2(t_0 + T) + \phi_3(t_0 + 2T) \quad (6.6)$$

where ϕ_i ($i = 1, 2, 3$) are the phases of the Raman fields on these three pulses, t_0 is the time at the start of the first $\pi/2$ pulse, and T is the time separation between interferometer pulses.

We only detect atoms in state $F = 3$ by the probe beam. When atoms are under free fall, the phases of three pulses are $\phi_1(t_0) = \omega t_0 + \phi_1^0$, $\phi_2(t_0 + T) = \omega(t_0 + T) - (\mathbf{k}_1 - \mathbf{k}_2) \cdot gT^2/2 + \phi_2^0$, and $\phi_3(t_0 + 2T) = \omega(t_0 + 2T) - (\mathbf{k}_1 - \mathbf{k}_2) \cdot g(2T)^2/2 + \phi_3^0$, where ω is the laser angular frequency, $|\mathbf{k}_1 - \mathbf{k}_2| \approx 2k_{eff}$ are the propagation constants of the Raman fields in the waveguide, g is the local gravity, and ϕ_i ($i = 1, 2, 3$) are the initial phases of the Raman fields. Therefore, the phase after closing the interferometer can be written as

$$\Delta\phi = -(\mathbf{k}_1 - \mathbf{k}_2) \cdot gT^2 + \phi_1^0 - 2\phi_2^0 + \phi_3^0 \quad (6.7)$$

where the first term is due to the Doppler shift of the free-falling atoms in the interferometer fields. Figure 6.3 shows the interference fringes by varying the laser phase ϕ_3^0 of the third pulse at different pulse separation times T , whereas the optical dipole trap is off during the interferometer time. The phase shift due to the Doppler shift of free-falling atoms becomes observable at longer T . The data are fitted by a sinusoidal function, and the contrast is defined by the

ratio of the difference and the sum of the maximum and minimum of fitted curves.

Figure 6.3: OD versus phase with different interferometer times T . The fringes are fitted to a sinusoidal function. Each data point is an average of 60 runs with the standard error.

In Fig. 6.4, we extract the phase shifts at different interferometer times by fitting the fringes to a sinusoidal function. We compare the values with our estimation. The local gravity in Singapore is obtained from International Gravity Formula 1980 (IFG80) as $g = 9.78 \text{ ms}^{-2}$, and the propagation constant $k_{eff} = 2\pi/(780 \times 10^{-9}) \text{ m}^{-1}$ is approximated as the wave vector in free space [76]. The parabolic increase of the phase shift over time indicates that the wavefront inside the hollow-core photonic crystal fiber can be used as a phase reference for coherent quantum system. The Fig. 6.5 shows the error of the gravity $\delta g/g$ as function of T^2 . The red curve is the fitting based on the Eq. 3.25.

Figure 6.4: Phase shift versus time T . Each data point is extracted from the fitting of a sinusoidal function as in Fig. 6.3 with the Raman beam direction along and opposite to the gravity. The zero point in the y axis corresponds to the phase at $T = 4 \mu\text{s}$. The curves are the fits with Eq. 6.3. The error bars are the uncertainty of the sinusoidal fits. The red circles are the phase shifts based on the estimation of Eq. 6.3.

Figure 6.5: The error of the gravity $\delta g/g$ as function of T^2 . The red curve is the fitting based on the Eq. 3.25.

Figure 6.6: Contrast versus time. A comparison of the decay time with optical dipole trap on until t_p and off during the interferometer sequence. The error bars are from the uncertainty of the sinusoidal fits.

We plot the contrast versus time T with dipole beam on and off during the interferometer time in Fig. 6.6. The decay of contrast with the dipole beam off is mainly due to the ballistic expansion of atoms in the fiber. When the dipole is on during the interferometer time, severe decay of contrast is observed. This is mainly due to the inhomogeneous broadening of differential ac stark shift $2\pi \times 3$ kHz between atomic states from the optical dipole beam.

Figure 6.7 shows interferometer fringes with $T = 100 \mu\text{s}$ at different probe times. The contrasts and phase shifts agree within one standard error. Using kinematic equation $gt_p^2/2$ from 25 to 35 ms, we operate the interferometer over 3-mm distance inside the fiber. This could be compared to the free-space parameter: The Rayleigh length $z_R = \pi W^2/\lambda$ of a Gaussian beam with waist

$W = 22$ mm is 1.95 mm, where l is the wavelength of the optical field.

Figure 6.7: OD versus phase at different probe times. Data are fitted to a sinusoidal function. The standard error is an average of 60 experimental runs.

6.3 Temperature of atoms in the atom interferometer

We also introduce timing asymmetry δT between three pulses to measure the velocity width of the atomic ensemble participating in the interferometric process as shown in Fig. 6.8 [91].

Figure 6.8: Velocity width measurement of atoms selected in the interferometer measurement. We use the sequence as shown in the inset. Each data point is from a sinusoidal fitting of interference fringe similar to Fig. 6.3 with $T = 4 \mu\text{s}$. The contrast decay is fitted to a Gaussian function. The error bars are from the uncertainty of the sinusoidal fitting.

We treat each atom as a coherent Gaussian wave packet. The wave packet can be expressed in space:

$$\psi(x) = \frac{1}{\sqrt{\sqrt{\pi}x_a}} \exp\left(-\frac{x^2}{2x_a^2}\right) \cdot \exp\left(i\frac{m}{h}v_c(x - x_0)\right) \quad (6.8)$$

where x_0 is the classical position of atom, v_c is the classical velocity of atom, and x_a is the coherence length of the wave packet. A set of momentum eigenstates can express the wave packet as

$$\psi(x) = \int_{-\infty}^{+\infty} dv \cdot \exp\left(-\frac{(v - v_c)^2}{2v_a^2}\right) \cdot \exp\left(i\frac{m}{h}v(x - x_0)\right) \quad (6.9)$$

where

$$v_a = \frac{\hbar}{mx_a}. \quad (6.10)$$

To find behavior of the ensemble, we should integrate over v_c and x_0 . We assume they obey the Gaussian distribution

$$f_v(v_c) \propto \exp\left(-\frac{v_c^2}{2v_t^2}\right) \quad (6.11)$$

v_t is defined by the rms thermal velocity

$$v_{rms} = \sqrt{v_t^2 + v_a^2/2}. \quad (6.12)$$

After the three Raman pulses (see Fig. 6.8), the wave packets of atom in two paths are

$$\psi_1(x) = \int_{-\infty}^{+\infty} dv \cdot \exp\left(-\frac{(v - v_c)^2}{2v_a^2}\right) \cdot \exp\left(i\frac{m}{h}v(x - x_0)\right) \quad (6.13)$$

$$\psi_2(x) = \int_{-\infty}^{+\infty} dv \cdot \exp\left(-\frac{(v - v_c)^2}{2v_a^2}\right) \cdot \exp\left(i\frac{m}{h}v(x - x_0)\right). \quad (6.14)$$

The probability to detect this atom is given

$$P_a = \int_{-\infty}^{+\infty} dx |\psi_1(x) + \psi_2(x)|^2. \quad (6.15)$$

The normalized atom number of detecting at this output port is

$$P = \int_{-\infty}^{+\infty} dv_c f_v(v_c) \cdot P_a = (1 + \chi \cos(\Delta\phi)) / 2 \quad (6.16)$$

where χ is the contrast given by

$$\chi = \exp\left(-\frac{(\delta T)^2}{(2x_a/v_r)^2}\right) \exp\left(-\frac{(\delta T)^2}{2/(k_{eff}v_t)^2}\right). \quad (6.17)$$

With $v_{rms} = \sqrt{v_t^2 + v_a^2/2}$ and $v_r = \hbar k_{eff}/m$ is the two-photon recoil velocity,

$$\chi = \exp\left(-\frac{(\delta T)^2}{2/(k_{eff}v_{rms})^2}\right). \quad (6.18)$$

We plot the contrast c versus the offset time δT and fit the contrast to a function $\chi(\delta T) = A \exp(-(v_{rms}k_{eff})^2(\delta T)^2/2)$, where A is the contrast at $\delta T = 0$. By assigning $k_{eff} \approx 2k = 2\pi/(780 \times 10^{-9}) \text{ m}^{-1}$, the fitted curve shows $v_{rms} = 3.17(0.25) \text{ cms}^{-1}$, where v_{rms} is the root mean square velocity of the atomic ensemble in the axial direction. Defining the thermodynamic temperature $T_p = mv_{rms}^2/k_B$, we estimate that the axial temperature of our atoms for the interferometer experiment is about $10 \text{ } \mu\text{K}$, where m is the mass of ^{85}Rb and k_B is the Boltzmann constant.

Chapter 7

Transporting long-lived quantum spin coherence inside fiber

One of the major tasks in modern quantum science and technology is to fully exploit the potentialities of quantum systems, by guiding and transporting quantum superposition states (spin qubits) of massive particles (atoms, ions and molecules) over a controllable distance while preserving their coherence [12, 34, 36, 85]. The hollow-core fibers provide a genuine platform for such tasks. However, the quantum superposition states of atoms are fragile from the trapping beams inside the fiber. The long-lived coherence of quantum superposition states inside the fiber remains challenging.

For neutral atoms, the restoring force used for confining atoms is generally through the coupling to the gradient of optical or magnetic fields. Coherent guiding of Bose-Einstein condensate using magnetic trap has been demonstrated on the scale of few hundred micrometers [43], but the ability to guide and maintain coherence over longer distances is challenging. Although, optical conveyor belts in free space have transported a superposition state of atoms over 1 mm range, the Rayleigh range over which a laser beam stays focused limits the guiding distances [44]. For optical nano-fibres, atoms are trapped a few hundred nanometers away from the surface by the evanescent waves of the guided mode. The coherence time of stationary atoms has been demonstrated

to few milliseconds [49] and will be fundamentally limited by the inhomogeneous broadening from the surface potential. Here, the dephasing mechanism of quantum superposition states of the atoms and the transportation of a quantum superposition state trapped inside the hollow-core photonic crystal fiber are studied. Our demonstration shows the first guiding of a quantum superposition state over a centimeter distance and preserve the coherence over hundreds of milliseconds inside the optical waveguide, improved ten and hundred-fold over the existing guiding schemes [44, 49], respectively. The transportation distance is limited only by our design of the experimental setup and the length of the fiber we used. Our results suggest that the interface of atoms with photonic crystal fibers could provide an important step in the realization of matter-wave circuit for remote and large scale quantum sensing [43, 92–96].

7.1 Atom loading using optical lattice

Figure 7.1: (a) Experimental setup of the loading using moving optical lattice. HCF: hollow-core fiber. QWP: quarter-wave plate. PBS: polarizing beam splitter. NPBS: non-polarizing beam splitter. APD: avalanche photodiode. PM fiber: polarization maintaining fiber. DM: dichroic mirror. (b) Timing scales for atoms' position in the hollow-core fiber. The left scale is the 50 kHz moving lattice sequence. The right scale represent a 50 kHz moving lattice sequence followed by the free-falling sequence starting at $t_p = 250$ ms.

Fig. 7.1 shows the experimental setup. Atoms are prepared in the MOT with temperature of $10\ \mu\text{K}$ and atom number of $\sim 10^9$. The MOT is aligned 5 mm above the fiber. Atoms are released from the MOT at $t_p = 0$. A moving optical lattice is then switched on to transfer atoms to the optical dipole trap.

We consider two counter-propagating beams with frequency detuned by Δf along the x axis. The electric fields of two beam are

$$E_1 = E_0 \cos(kx - (2\pi(f + \Delta f)t)) \quad (7.1)$$

and

$$E_2 = E_0 \cos(kx - 2\pi ft) \quad (7.2)$$

where $k = 2\pi/\lambda$ is the wave number and $\lambda = 875$ nm is the wavelength of laser beams. The amplitude of the oscillations of the optical lattice along x axis is given by

$$A(x) = 2E_0 \cos[2k(x - \frac{\pi\Delta f}{k}t)]. \quad (7.3)$$

The lattice will move in the direction of the laser with higher-frequency. The speed of moving lattice can be written as

$$v = \frac{\pi\Delta f}{k} \quad (7.4)$$

with $k = 2\pi/\lambda$, the velocity of the lattice is

$$v = \frac{\lambda\Delta f}{2}. \quad (7.5)$$

Atoms are loaded into the fiber by the moving lattice. The push beam resonant on the $F = 3$ to $F' = 4$ transition is used to push atoms outside the fiber away. When atoms are loaded into the fiber, the frequency detuning Δf is switched to zero forming a stationary lattice to hold the atoms. Figure 7.2 shows the fluorescence images of atoms in the moving optical lattice. The insert graph shows the distribution of atoms in the radial direction at 135 ms, it is a Gaussian distribution. After 155 ms, atoms are loaded into the fiber completely. This result gives a direct evidence to prove that atoms are transported into the hollow-core fiber using the moving lattice. The data shown in Fig. 7.3 is the OD as a function of traveling time with frequency detuning $\Delta f = 50, 100, 150, 200$, and 300 kHz, respectively.

Figure 7.2: The fluorescence images of atoms during the loading process. The insert graph shows the atoms distribution in the radial direction is a Gaussian distribution.

Figure 7.3: OD as a function of traveling time with frequency detuning $\Delta f = 50, 100, 150, 200,$ and 300 kHz.

7.2 Lifetime of atoms inside the fiber

Figure 7.4: Lifetime of atoms in the stationary lattice at different positions inside the fiber.

The moving lattice is used to transfer atoms to different positions inside the fiber. A stationary lattice is then formed to hold atoms at the same position. The data is shown in Fig. 7.4. The frequency detuning of the moving lattice is 50 kHz. The stationary lattice is applied at $t_p = 500, 700$, and 900 ms for lifetime measurements. The probe beam resonant to $F = 3$ to $F' = 4$ transition is sent for OD measurements. The data shows the decay of OD as function of time. The curves are fitted to an exponential decay function. The $1/e$ decay time of the measurements is 1.87, 2.59, and 2.33 s, respectively. Our result suggests that there is no significant heating loss of atoms inside the hollow-core fiber. The result of lifetimes also indicates that the background pressure

inside the hollow-core fiber is better than 1×10^{-8} torr [28]. The guiding environment is clean enough for our experiments.

7.3 Temperature measurements of atoms inside the fiber

Figure 7.5: The temperature of atoms inside the fiber using moving lattice for the transportation.

The temperature of atoms inside the fiber is measured by the method used in section 6.3. When atoms are loaded into the fiber by using the moving lattice with frequency detuning 50 kHz, the dipole beams are switched off and the probe beam resonant to $F = 3$ to $F' = 4$ transition is sent after some delay time. The result is shown in Fig. 7.5. We measure the temperature at different transportation time t_p . Then we plot the temperature-time cure. We found that the temperature is not increased significantly during the trans-

portation, so there is no extra heating of atoms by using the moving lattice for the transportation.

The temperature of atoms inside fiber is about $\sim 80 \mu\text{K}$. Compared to the initial temperature of atoms prepared in the MOT of $\sim 10 \mu\text{K}$, the increase of the temperature is due to loading process of atoms into the optical dipole trap at the tip of the fiber. At the tip of fiber, a partial dipole beam is reflected by the tip. A residual standing wave is created by the reflection which modulate the moving lattice potential as a source of parametric heating on atoms. After loading atoms into the fiber, there is no residual standing wave and atoms are not heated inside the moving lattice. In the future, the heating mechanism can be studied in details.

7.4 Coherence time of atoms inside fiber

Figure 7.6: Time sequences. The external magnetic fields B_x , B_y , and B_z are for cancelling the ambient magnetic field and defining quantization axis.

The Ramsey and spin echo techniques are employed to study the dephasing mechanism of atoms inside the fiber. We send two microwave pulses ($\pi/2 -$

$T - \pi/2$) to form the Ramsey interferometer and three pulses ($\pi/2 - T/2 - \pi - T/2 - \pi/2$) to form the spin-echo interferometer, where T is the pulse separation time. They are used to study the coherence of the ground state $F = 2, m = 0$ and $F' = 3, m = 0$ of ^{85}Rb . The frequency of the microwave is 3.04 GHz resonant to the transition of two hyperfine ground states. The timing sequence is shown in Fig. 7.6. After loading atoms into the fiber, the push beam is sent to push atoms outside the fiber away.

Figure 7.7: Zeeman spectroscopy of the atoms inside the fiber after the optical pumping

Before implementing the interferometer, an optical pumping is used to bring atoms to the $F = 2, m = 0$ state. During the optical pumping, a magnetic field is applied in x direction to define the quantization axis. After the optical pumping, the magnetic field is changed to the z direction for the interferometer

implementation. A Zeeman spectroscopy of the atoms inside fiber after the optical pumping is shown in Fig. 7.7. The efficiency of the optical pumping is about 71%. The optical pumping increases the population of atoms used in the interferometer.

Figure 7.8: Coherence time measurement of stationary atoms inside the hollow-core fiber. The Ramsey and spin echo techniques are employed.

Due to the non-uniform transverse profile of the trapping beam inside the fiber, the inhomogeneous differential ac Stark shift of the transition is present on the two hyperfine ground states. Due to the inhomogeneous broadening introduced by the dipole beams which limits the coherence time to few milliseconds. The results are shown in Fig. 7.8. The trapping potential of the stationary lattice is $537 \mu\text{K}$. The blue data is the Ramsey signal, and the coherence time of the atoms is about 3 ms. We then apply the spin-echo technique to remove the dephasing due to the inhomogeneous broadening of the

transition. In principle, the spin-echo technique can reverse the sign of the phase shift by adding a π pulse between the two $\pi/2$ pulses to remove the dephasing [97]. By this method, the coherence time is increased to ~ 37 ms (red data points).

To model the action of the spin-echo pulse sequence, the optical Bloch equations [98] is introduced

$$\dot{\mathbf{u}} = -\mathbf{\Omega} \times \mathbf{u}. \quad (7.6)$$

Here, $\mathbf{\Omega} \equiv (\Omega, 0, \delta)$ is the torque vector, and $\mathbf{u} \equiv (u, v, w)$ is the Bloch vector. Ω is the Rabi frequency, and δ is the detuning of the microwave pulses. After the first $\pi/2$ pulse at $t = 0$, the Bloch vector starts at $\mathbf{u}_0 \equiv (0, -1, 0)$. The inhomogeneous dephasing will make the Bloch vectors to rotate around the w axis by slightly different frequencies. After a π pulse at time T , the Bloch vectors rotate around the u axis by 180° and a complete rephasing is induced at time $2T$ in the state $\mathbf{u}(2T) \equiv (0, 1, 0)$.

We use the rotation matrices acting on the Bloch vector to express the solutions of the optical Bloch equations. The spin-echo sequence then reads

$$\mathbf{u}_{echo}(t) = \Theta_{\pi/2} \cdot \Phi_{free}(\delta, t - T) \cdot \Theta_{\pi} \cdot \Phi_{free}(\delta, T) \cdot \Theta_{\pi/2} \cdot \mathbf{u}_0. \quad (7.7)$$

The action of a $\pi/2$ pulse is described by the matrices

$$\Theta_{\pi/2} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & -1 & 0 \end{pmatrix} \quad (7.8)$$

and the action of a π pulse is described by

$$\Theta_{\pi} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & -1 \end{pmatrix}. \quad (7.9)$$

The free precession around the w axis during a time t is given by

$$\Phi_{free}(\delta, t) = \begin{pmatrix} \cos(\phi(\delta, t)) & \sin(\phi(\delta, t)) & 0 \\ -\sin(\phi(\delta, t)) & \cos(\phi(\delta, t)) & 0 \\ 0 & 0 & 1 \end{pmatrix}$$

with $\phi(t) = \int_0^t \delta(t') dt'$ is the total free precession angle during the free evolution of the Bloch vector.

We assume a time-varying detuning $\delta(t)$, the free precession angle before and after the π pulse are

$$\phi_1 = \int_0^T \delta(t) dt \quad (7.10)$$

and

$$\phi_2 = \int_T^{2T} \delta(t) dt. \quad (7.11)$$

The phase difference is

$$\Delta\delta = \frac{\phi_2 - \phi_1}{T}. \quad (7.12)$$

At time $2T$, the result of Eq. (7.2) is expressed in

$$w_{echo}(\Delta\delta, 2T) = -\cos(\Delta\delta T). \quad (7.13)$$

If $\Delta\delta$ is a Gaussian distribution with mean $\overline{\Delta\delta} = 0$ and variance $\sigma(T)^2$, the average w component of the Bloch vector is

$$w_{echo}(2T) = \int_{-\infty}^{\infty} -\cos(\Delta\delta T) \frac{1}{\sigma(T)\sqrt{2\pi}} \exp\left[-\frac{(\Delta\delta)^2}{2\sigma(T)^2}\right] d\Delta\delta' = \exp\left[-\frac{1}{2}T^2\sigma(T)^2\right]. \quad (7.14)$$

Thus, the visibility of the spin-echo V is

$$V(2T) = V_0 \exp\left[-\frac{1}{2}T^2\sigma(T)^2\right]. \quad (7.15)$$

The $1/e$ decay time of the visibility is the homogeneous dephasing time T'

$$T' = \frac{\sqrt{2}}{\sigma(T)}. \quad (7.16)$$

To increase the coherence time, we use the vector light shift to cancel the differential scalar shift. The polarization of the dipole beams is adjusted by a quarter wave-plate to tune the vector light shift. The applied magnetic field is about 2.2 G. Although this method has been pointed out by the previous studies that the perfect cancellation does not exist for the hyperfine ground states of alkaline-metal atoms [99, 100], we still follow the method to improve the coherence time to 60.1(1) ms after the cancellation.

The differential light shift of two hyperfine ground states in presence of a magnetic field can be calculated as

$$\delta\nu(B, U) = \beta_1 U + \beta_2 BU + \beta_4 U^2 \quad (7.17)$$

where U is the trapping potential, B is the external magnetic field, and β_1 , β_2 , and β_4 are the coefficients of the light shift. β_1 is given by

$$\beta_1 = [\alpha_{F=3}(\omega) - \alpha_{F=2}(\omega)]/\alpha_{5s}(\omega) \quad (7.18)$$

with $\alpha_F(\omega)$ are the hyperfine-interaction-mediated polarizability of the hyperfine ground states and $\alpha_{5s}(\omega)$ is the ground state polarizability. β_2 is expressed as

$$\beta_2 = -4A\mu_B\alpha^v(\omega)/[hf\alpha_{5s}(\omega)] \quad (7.19)$$

where μ_B is the Bohr magneton, $\alpha^v(\omega)$ is the vector polarizability of the ground state, A is the polarization of the trapping beams, f is the frequency of hyperfine splitting, and h is the Planck constant. β_4 can be expressed as

$$\beta_4 = (A^2/2f)[\alpha^v(\omega)/\alpha_{5s}(\omega)]^2 \quad (7.20)$$

In our experiment, we calculate and use $\beta_1 = 7.86 \times 10^{-5}$, $\beta_2 = -1.13 \times 10^{-4} \text{ G}^{-1}$, and $\beta_4 = 2.48 \times 10^{-12} \text{ Hz}^{-1}$. By adjusting the polarization of the trapping beams, the differential light shift of two hyperfine ground states ($\delta\nu(B, U)$) in presence of a magnetic field can be partially cancelled.

After optimizing the cancellation, we also measure the coherence time of atoms at different positions inside the hollow-core fiber. Atoms are transferred by the 50 kHz moving lattice and hold by the stationary lattice after traveling for 250 ms, 550 ms, and 750 ms, respectively. The traveling distances are about 5.5, 12.0, 16.4 mm, respectively (start from the center of the MOT). As shown in Fig. 7.9, the coherence time of atoms are 106(2), 108(2), and 106(2). The result shows that the coherence time is consistent in different positions along the hollow-core fiber.

Figure 7.9: Coherence times of atoms at different positions inside the fiber.

Figure 7.10: Spin-echo measurement with light shift cancellation of two different trapping potential 279 μ K (squares) and 179 μ K (circles).

We also compare the spin-echo signal at two different trapping potentials. Figure 7.10 shows the data of coherence time measurement at trapping potential 279 μ K to be 101(2) ms. We then reduce the trapping potential to 179 μ K, and observe the improvement of the coherence time to 252(8) ms. The single photon scattering rates at the peak trapping potentials are 2.88 s^{-1} and 6.79 s^{-1} respectively, which are not the limitation of the coherence time. The source of decoherence on atoms is the residual inhomogeneous differential ac Stark shift.

7.5 Transportation of quantum superposition

To transport the coherent superposition state of atoms in the hollow-core fiber, a $\pi/2$ microwave pulse is applied at $t_p = 250$ ms to prepare atoms in a superposition state. The relative phase of the two states evolves when atoms are in motion along the fiber. Finally, the interference is read out either by Ramsey or spin-echo sequences. Figure 7.11 shows the coherence time measurement of atoms in a 50 kHz moving lattice with different pulse separation time T with (squares) and without (circles) light shift cancellation. With the light shift cancellation, the distance of the transportation is about 1.75 mm after traveling for 80 ms.

Figure 7.11: Coherence time of moving atoms inside the fiber.

We transport atoms in the moving lattice for 250 ms and then switch off the bottom dipole beam at $t_p = 250$ ms to transfer atoms to the free-falling frame guided only by the top dipole beam. In Fig. 7.12, we compare the results

of different measurement conditions when atoms are under free fall. Without the cancellation, the coherence time of the spin-echo sequence greatly exceed the Ramsey sequence measurement while with the cancellation. With the light shift cancellation, the coherence time using spin-echo technique didn't outperform the coherence time of Ramsey method. With the spin-echo technique and the light shift cancellation, the coherence time of moving atoms is about 45 ms, which corresponds to a transportation distance of 1 cm.

Figure 7.12: Coherence time of free-falling atoms inside the fiber.

The source of decoherence on the moving atoms is mainly due to the magnetic field gradient. To estimate the magnetic field gradient of our setup, we measure the linewidth of magnetic field sensitive transition $F = 2, m = 1$ to $F' = 3, m = 2$ using microwave spectroscopy and compare it with the magnetic field insensitive transition $F = 2, m = 1$ to $F' = 3, m = 0$. Figure 7.13 shows the measurement of the transition linewidth of 11.6 kHz and 1.55 kHz for magnetic field sensitive and insensitive transition respectively. Assume the broadening on the insensitive transition is mainly due to the power broaden-

ing of the microwave, we fit the sensitive transition of a convolution of double Lorentzian functions, one is from the power broadening and the other is due to magnetic field gradient. The fit result shows a gradient of 7.1 mG over the sample size of about 5 mm. This gradient translates into the broadening of about 25 Hz on the $F = 2, m = 1$ to $F' = 3, m = 0$ transition due to the second order Zeeman shift at the magnetic field of 2.2 G.

Figure 7.13: Zeeman spectroscopy of atoms trapped in a linear polarized stationary optical lattice. An external 2.2 G magnetic field is applied to break the Zeeman degeneracy. Top figure is the transition from $F = 2, m = 1$ to $F' = 3, m = 0$ and the bottom figure is $F = 2, m = 1$ to $F' = 3, m = 2$.

Chapter 8

Conclusions

8.1 Summary

In this thesis, I studied optically trapped atoms inside a hollow-core photonic crystal fiber and the use of optical waveguide mode inside the fiber as beam splitter and mirror pulses to manipulate atoms to form an inertia-sensitive atom interferometer which can be used to measure the acceleration due to gravity. The distance that atoms can propagate is not limited by the diffraction of the light of the interferometer pulses. Moreover, the dephasing mechanism of quantum superposition states of atoms and the transportation of quantum superposition inside the hollow-core photonic crystal fiber are studied. Coherent interactions between the waveguide mode and quantum systems could lead to constructing mobile and miniature quantum hybrid platforms for precision measurement and quantum sensing. These demonstrations show the first guiding of a quantum superposition state of trapped atoms over a centimeter distance inside the optical waveguide and preserve the coherence over hundreds of milliseconds. Our results suggest that atom-light-interactions in a hollow-core fiber show great promise to the realization of fiber based matter-wave circuits, quantum sensors, and quantum memories.

8.2 Outlooks

Recently, the degenerate Raman sideband cooling of ^{85}Rb atoms in a two-dimensional optical lattice have been demonstrated [101]. It shows the cooling of spin-polarized ^{85}Rb atoms to the recoil temperature. We will follow this work to implement one-dimensional Raman sideband cooling along the fiber axis to cool down the atoms inside the fiber. Atom numbers and the contrast of our interferometer can be improved by lowering the temperature of the atomic ensemble. The temperature in the radial direction can be cooled down through collisions or coupling between axial and radial motion in the optical lattice. The two counter-propagating dipole beams are used to implement the degenerate two-photon Raman transition for the interchange of vibrational energy and Zeeman energy. We will employ a D_1 line for optical pumping to avoid coupling to other excited states. The σ^+ -polarized and the π -polarized optical pumping beams will be used to perform the cooling processing. The heating due to the dipole intensity noise can be minimized by implementing an intensity stabilization feedback system. With 10^4 atoms and $T = 1$ ms coherence time, our interferometer can have a relative sensitivity of the propagation constant $\Delta k/k = 6 \times 10^{-5}$ with a distance of the separation of atomic wave packets $2v_r T \sim 10 \mu\text{m}$ averaging over the size of the atomic ensemble, where v_r is the photon recoil velocity of the laser used here.

The fiber is not exactly in vertical direction. This is one of the systematic error. Our measurements get gravity as $g = 9.75 \pm 0.3 \text{ m/s}^2$. The local gravity in Singapore is obtained from International Gravity Formula 1980 (IGF80) as $g = 9.78 \text{ m/s}^2$. From one error bar of our result, I estimate the tilt angle is $\sim 4.5^\circ$. In the future, we will put the fiber on a rotation stage. By changing the rotation angle of the fiber, we can study the systematic error from the alignment of the fiber.

With the combination of the photonic system, low power and small volume of our interferometer shows prospects of compactness and versatility. For example, in Newton's constant G measurement [102] and the proposed gravitational Aharonov-Bohm experiment [103], our interferometer could permit a

short-distance arrangement between source masses and atoms to enhance the signal. In the study of Rosi et al. [102], with the same amount of source masses and interferometer performance, our interferometer could increase the phase shift by a factor of 7. More importantly, although the size and the position of the atomic cloud are the leading systematic effects at the level of hundreds of micrometers, our interferometer can have an improvement by two orders of magnitude with 10% uncertainty of the fiber position. The precision (or uncertainty) of any atom interferometers is fundamentally limited by the atom shot noise. The precision of the work of Rosi et al. [102] is 0.012 rad with 2×10^5 atoms used, a factor of 5 away from atom shot noise. Although our fiber atom interferometer only uses about 10^4 atoms, orders of magnitude improvement can be made in the future by cooling atoms to lower temperature and optical pumping to the magnetic field-insensitive state. By preparing optically thick atomic medium in the fiber, spin-squeezing techniques can be applied to improve the sensitivity of the interferometer [104, 105]. Our experiment demonstrated here will also initiate a campaign of studying the systematic effects of atom interferometers in the hollow-core fiber such as the propagation constant and imperfection of the fiber mode structure. Our realization of an atom interferometer inside a hollow-core photonic crystal fiber indicates coherent interaction of the optical waveguide mode and quantum system, and could open up novel applications for precision measurement and quantum sensing.

The coherence time of stationary atoms in our experiment can be improved by exciting ring-shaped higher order mode of the fiber with blue detuned dipole laser frequency or using alkaline-earth metal atoms which have magic-wavelengths for narrow-linewidth optical transitions. The distance of the optical waveguide enhanced quantum state transportation demonstrated in this work have already been an order of magnitude longer than the free space schemes. We expect the coherence time can be improved by another order of magnitude by wrapping the fiber with high permeability but soft materials to shield the magnetic field gradient and further minimizing the residual ac Stark shift with colder atoms and shallower trapping potential. This will translate into more than 1 meter transportation of the quantum state in the free falling

case. Considering a fiber with 5 cm bending radius at 3 dB bending loss, our work can be used to construct a matter-wave circuit with 30 cm perimeter for quantum sensing purposes. A fiber of miniaturized quantum devices could be envisioned.

Quantum memories are regarded as one of the fundamental building blocks of quantum computation [106, 107]. A long-standing goal to realize scalable quantum information processing is to achieve an efficient and long-lived quantum memory. So far, either high-efficient but short-lived [108] or long-lived but low-efficient quantum memories [109] have been demonstrated. Long-lived storage with high efficiency remains an outstanding challenge. Quantum memories for the storage and retrieval of quantum information are extremely sensitive to environmental influences, which limits their storage times. In the stationary atoms in our experiment, we have achieved the long coherence time. In the future, we will try to perform the quantum memory using atoms inside fiber to achieve long storage time. Due to the strong confinement of atoms inside the fiber, we have achieved high optical depth of atoms inside fiber. The higher optical depth leads to high storage efficiency in quantum memory. We expect to achieve long storage time with high efficiency using the new scheme. This realization will provide an essential tool towards scalable quantum information processing and communications applications.

Free-space atomic systems are mature technologies in research laboratories. How to miniaturize these table-top apparatus for real-world applications is still an open question. The main obstacle comes from the diffraction of light since most of the quantum computing, communications, and sensing applications based on atoms relies on atom-light interactions. Coupling atoms to diffraction-free waveguide systems while preserving their spin coherence is one possible route. Our results could pave the way for the realization of miniaturized atomic systems.

Appendix

Appendix A: Characteristic data of ^{85}Rb

Quantity	Symbol	Value (SI)
Melting point	T_M	39.30°C
Atomic mass	m	$1.409\,993\,199(70) \times 10^{-25}\text{kg}$
Frequency	ω_0	$2\pi \cdot 384.230\,406\,373(14)\,\text{THz}$
Wavelength (Vacuum)	λ	$780.241\,368\,271(27)\,\text{nm}$
Lifetime	τ	$26.2348(77)\,\text{ns}$
Natural linewidth	Γ	$2\pi \cdot 6.0666(18)\,\text{MHz}$
Hyperfine splitting	ω_{HF}	$3.035\,732\,439\,0(60)\,\text{GHz}$
Clock transition Zeeman shift	$\Delta\omega_{clock}/B^2$	$2\pi \cdot 1293.98\,\text{Hz/G}^2$
Doppler temperature	T_D	$145.57\,\mu\text{K}$
Doppler velocity	v_D	$9.52\,\text{cm/s}$
Recoil temperature	T_r	$370.47\,\text{nK}$
Recoil velocity	v_r	$6.0230\,\text{mm/s}$
Saturation intensity	I_s	$3.895\,09(81)\,\text{mW/cm}^2$

Figure 8.1: Physical properties of ^{85}Rb and optical properties of D_2 line transition ($5^2S_{1/2} \rightarrow 5^2P_{3/2}$).

Quantity	Symbol	Value (SI)
Speed of light	c	$2.997\,924\,58 \times 10^8$ m/s (exact)
Planck's constant	$\hbar$	$1.054\,571\,628(53) \times 10^{-34}$ J·s $6.582\,118\,99(16) \times 10^{-16}$ eV·s
Boltzmann's constant	k_B	$1.380\,6504(24) \times 10^{-23}$ J/K
Bohr magneton	μ_B	$9.274\,009\,15(23) \times 10^{-24}$ J/T $\hbar \cdot 1.399\,624\,604(35)$ MHz/G

Figure 8.2: Fundamental physical constants used in the thesis.

Appendix B: Fiber preparation

The first version of fiber we used is the mode number HC-800B from Thorlabs. The mode field diameter is $5.5 \pm 1 \mu\text{m}$. We used a lens with focus length of 18 mm to couple light into the fiber. The lens was placed inside the vacuum chamber, and the position of the lens was fixed. The coupling is hard to implement and efficiency is about 40%. The profile of the output beam is nearly Gaussian (see Fig. 8.3).

Figure 8.3: (a) Microscope image of the cross section of the hollow-core fiber (HC-800B). (b) Typical near field intensity profile of the beam through the fiber.

Figure 8.4: (a) The microscope image of the blocked fiber. (b) The wire-connector used for the rubidium dispenser.

Figure 8.5: (a) The microscope image of the new fiber. (b) Side view of the fibers placed on the fiber holder.

The hollow-core fiber used in this work was bought from GLOphotonics of model number *PMC-C-TiSa_Er-7C*. The first fiber we used was blocked by some liquid, which is shown in Fig. 8.4 (a), and the tip was destroyed by the laser beams coupled to the fiber. The liquid is oily and is not from the process of manufacture. We blame it to the wire connector used for the rubidium dispenser (see Fig. 8.4(b)).

The rubidium dispenser provides the source of rubidium atoms for the experiment. By running a DC current to heat the dispenser through ohmic heating, rubidium atoms evaporate to the vacuum chamber.. Due to the high temperature ~ 500 °C during ohmic heating, the coating of the wire connected to the dispenser burned and the generated debris scattered onto the fiber tips. Therefore, we changed the rubidium dispenser by a rubidium ampoule. We also changed the fiber destroyed. Three new fibers are then placed on the fiber holder. The microscope image of one of the new fibers is shown in Fig. 8.5.

Appendix C: Magnetic coil and water cooling system

A pair of anti-Helmholtz coils generates a spherical-quadrupole-symmetry magnetic field in all three directions. Each coil has 80 turns coiling by squared copper wire. The distribution of the magnetic field along the main coil axis is shown in Fig. 8.6. With current 17 A, the gradient of the magnetic field near the center is about 12 G/cm. It is enough for the MOT building.

A high current runs through the magnetic coils during the laser cooling process. With the current of 17 A and voltage of 6 V, the coils dissipate 102 W of power. The water cooling is required for the high power dissipation of the magnetic coils. We design a coil mount with empty space in the middle through which coils is placed in and cooling water is run. Before putting the coils into the mount, we paint the coils using a thermal conducting resin to make coils isolate from the water. Fig. 8.7(a) shows the coils after painting. The resin is made of STYCAST 2850MT and a catalyst. After painting, the coils are placed into the mount, and a cover is welded to the mount. The design of the coil mount is shown in Fig. 8.7(b).

Figure 8.7: (a) Two magnetic coils painting with a resin. (b) The design of the coil mount and the cover.

Figure 8.6: Magnetic field distribution of the anti-Helmholtz coils along the main coil axis.

Appendix D: Number and temperature of atoms inside fiber

Optical Depth

Normalized atom distribution

$$n(r, t) = \frac{1}{\pi r_0^2(t)} e^{-\frac{r^2}{r_0^2(t)}} \quad (8.1)$$

where $r_0^2(t) = r_0^2 + v_0^2 t^2$. The Gaussian beam distribution goes like

$$I(r) \sim e^{-\frac{r^2}{w^2/2}} \quad (8.2)$$

So,

$$\begin{aligned} \text{OD} &= \text{OD}_0 \int_0^R n(r, t) e^{-\frac{r^2}{w^2/2}} \cdot 2\pi r dr \\ &= \text{OD}_0 \frac{W^2/2}{r_0^2 + v_0^2 t^2 + W^2/2} \left(1 - e^{-\frac{R^2}{r_0^2 + v_0^2 t^2} - \frac{R^2}{W^2/2}} \right) \end{aligned} \quad (8.3)$$

where OD_0 is the OD at $t = 0$ with infinitesimal thin cloud size.

Equipartition theorem

In equilibrium, equipartition theorem tells us that

$$\begin{aligned} \left\langle \frac{1}{2} m v_0^2 \right\rangle &= k_B T \\ \Rightarrow \langle v_0^2 \rangle &= \frac{2k_B T}{m} \end{aligned} \quad (8.4)$$

and

$$\left\langle \frac{1}{2} k r_0^2 \right\rangle = k_B T \quad (8.5)$$

where k is the “spring constant” in simple harmonic oscillator (approximate) from Gaussian dipole trap (W =beam waist)

$$\begin{aligned}
 U &= -U_0 e^{-\frac{2r^2}{W^2}} \\
 &= -U_0 + 2\frac{U_0 r^2}{W^2} + O(r^4) \\
 &\equiv -U_0 + \frac{1}{2}kr^2
 \end{aligned} \tag{8.6}$$

where $k = \frac{4U_0}{W^2}$. So continue,

$$\Rightarrow \langle r_0^2 \rangle = \frac{k_B T}{2U_0} W^2. \tag{8.7}$$

Optical Depth vs Temperature

Combining everything,

$$\begin{aligned}
 \text{OD} &= \text{OD}_0 \frac{W^2/2}{r_0^2 + v_0^2 t^2 + W^2/2} \left(1 - e^{-\frac{R^2}{r_0^2 + v_0^2 t^2} - \frac{R^2}{W^2/2}} \right) \\
 &= \text{OD}_0 \frac{W^2/2}{\frac{k_B T}{2U_0} W^2 + \frac{2k_B T}{m} t^2 + W^2/2} \left(1 - e^{-\frac{k_B T}{2U_0} \frac{R^2}{W^2} - \frac{2k_B T}{m} \frac{R^2}{t^2} - \frac{R^2}{W^2/2}} \right) \\
 &= \frac{\text{OD}_0}{\frac{k_B T}{U_0} + \frac{4k_B T}{mW^2} t^2 + 1} \left(1 - e^{-\frac{k_B T}{2U_0} \frac{R^2}{W^2} - \frac{2k_B T}{m} \frac{R^2}{t^2} - \frac{R^2}{W^2/2}} \right) \\
 &= \frac{\text{OD}_0}{\frac{k_B T}{U_0} + \frac{4k_B T}{mW^2} t^2 + 1} \left(1 - e^{-\frac{R^2}{W^2/2} \left(\frac{k_B T}{2U_0} \frac{W^2/2}{W^2} + \frac{2k_B T}{m} t^2 + 1 \right)} \right) \\
 &= \frac{\text{OD}_0}{\frac{k_B T}{U_0} + \frac{4k_B T}{mW^2} t^2 + 1} \left(1 - e^{-\frac{R^2}{W^2/2} \left(\frac{k_B T}{U_0} + \frac{4k_B T}{mW^2} t^2 + 1 \right)} \right) \\
 &\equiv \frac{\text{OD}_0}{AT + BTt^2 + 1} \left(1 - e^{-C \left(\frac{1}{AT + BTt^2} + 1 \right)} \right).
 \end{aligned} \tag{8.8}$$

Here I use $U_0 = U_0 = 0.53$ mK, $W = 22$ μm , $R = 31.5$ μm , $m = 1.4431606 \times 10^{-25}$ kg, which means that $A = 0.001909091$ μK^{-1} , $B = 0.790648$ $\mu\text{K}^{-1}\text{ms}^{-2}$,

$C = 4.10021$.

Bibliography

- [1] Savas Dimopoulos, Peter W. Graham, Jason M. Hogan, and Mark A. Kasevich. Testing general relativity with atom interferometry. *Physical Review Letters*, 98(11), mar 2007.
- [2] Rym Bouchendira, Pierre Cladé, Saïda Guellati-Khélifa, François Nez, and François Biraben. New determination of the fine structure constant and test of the quantum electrodynamics. *Physical Review Letters*, 106(8), feb 2011.
- [3] Richard H. Parker, Chenghui Yu, Weicheng Zhong, Brian Estey, and Holger Müller. Measurement of the fine-structure constant as a test of the standard model. *Science*, 360(6385):191–195, apr 2018.
- [4] J. B. Fixler, G. T. Foster, J. M. McGuirk, and M. A. Kasevich. Atom interferometer measurement of the newtonian constant of gravity. *Science*, 315(5808):74–77, jan 2007.
- [5] G. Lamporesi, A. Bertoldi, L. Cacciapuoti, M. Prevedelli, and G. M. Tino. Determination of the newtonian gravitational constant using atom interferometry. *Physical Review Letters*, 100(5), feb 2008.
- [6] M. Kasevich and S. Chu. Measurement of the gravitational acceleration of an atom with a light-pulse atom interferometer. *Applied Physics B Photophysics and Laser Chemistry*, 54(5):321–332, may 1992.
- [7] A. Peters, K. Y. Chung, B. Young, J. Hensley, and S. Chu. Precision atom interferometry. *Philosophical Transactions of the Royal Society A: Mathematical, Physical and Engineering Sciences*, 355(1733):2223–2233, dec 1997.
- [8] R. Geiger, V. Ménoret, G. Stern, N. Zahzam, P. Cheinet, B. Battelier, A. Villing, F. Moron, M. Lours, Y. Bidel, A. Bresson, A. Landragin, and P. Bouyer. Detecting inertial effects with airborne matter-wave interferometry. *Nature Communications*, 2(1), sep 2011.

-
- [9] M. Snadden, J. McGuirk, P. Bouyer, K. Haritos, and M. Kasevich. Measurement of the earth's gravity gradient with an atom interferometer-based gravity gradiometer. *Physical Review Letters*, 81(5):971–974, aug 1998.
 - [10] G. Rosi, L. Cacciapuoti, F. Sorrentino, M. Menchetti, M. Prevedelli, and G.M. Tino. Measurement of the gravity-field curvature by atom interferometry. *Physical Review Letters*, 114(1), jan 2015.
 - [11] T. L. Gustavson, P. Bouyer, and M. A. Kasevich. Precision rotation measurements with an atom interferometer gyroscope. *Physical Review Letters*, 78(11):2046–2049, mar 1997.
 - [12] Alexander D. Cronin, Jörg Schmiedmayer, and David E. Pritchard. Optics and interferometry with atoms and molecules. *Rev. Mod. Phys.*, 81:1051–1129, Jul 2009.
 - [13] Achim Peters, Keng Yeow Chung, and Steven Chu. Measurement of gravitational acceleration by dropping atoms. *Nature*, 400:849, August 1999.
 - [14] Holger Müller, Sheng wey Chiow, Quan Long, Sven Herrmann, and Steven Chu. Atom interferometry with up to 24-photon-momentum-transfer beam splitters. *Physical Review Letters*, 100(18), may 2008.
 - [15] Shau-Yu Lan, Pei-Chen Kuan, Brian Estey, Philipp Haslinger, and Holger Müller. Influence of the coriolis force in atom interferometry. *Physical Review Letters*, 108(9), feb 2012.
 - [16] Ying-Ju Wang, Dana Z. Anderson, Victor M. Bright, Eric A. Cornell, Quentin Diot, Tetsuo Kishimoto, Mara Prentiss, R. A. Saravanan, Stephen R. Segal, and Saijun Wu. Atom michelson interferometer on a chip using a bose-einstein condensate. *Physical Review Letters*, 94(9), mar 2005.
 - [17] T. Schumm, S. Hofferberth, L. M. Andersson, S. Wildermuth, S. Groth, I. Bar-Joseph, J. Schmiedmayer, and P. Krüger. Matter-wave interferometry in a double well on an atom chip. *Nature Physics*, 1(1):57–62, sep 2005.
 - [18] M. J. Renn, D. Montgomery, O. Vdovin, D. Z. Anderson, C. E. Wieman, and E. A. Cornell. Laser-guided atoms in hollow-core optical fibers. *Physical Review Letters*, 75(18):3253–3256, oct 1995.
 - [19] R. F. Cregan. Single-mode photonic band gap guidance of light in air. *Science*, 285(5433):1537–1539, sep 1999.

-
- [20] P. J. Roberts, F. Couny, H. Sabert, B. J. Mangan, D. P. Williams, L. Farr, M. W. Mason, A. Tomlinson, T. A Birks, J. C. Knight, and P. St. J. Russell. Ultimate low loss of hollow-core photonic crystal fibres. *Optics Express*, 13(1):236, 2005.
- [21] Y. Y. Wang, N. V. Wheeler, F. Couny, P. J. Roberts, and F. Benabid. Low loss broadband transmission in hypocycloid-core kagome hollow-core photonic crystal fiber. *Optics Letters*, 36(5):669, feb 2011.
- [22] S Vorrath, S A Möller, P Windpassinger, K Bongs, and K Sengstock. Efficient guiding of cold atoms through a photonic band gap fiber. *New Journal of Physics*, 12(12):123015, dec 2010.
- [23] M. Bajcsy, S. Hofferberth, T. Peyronel, V. Balic, Q. Liang, A. S. Zibrov, V. Vuletic, and M. D. Lukin. Laser-cooled atoms inside a hollow-core photonic-crystal fiber. *Physical Review A*, 83(6), jun 2011.
- [24] M. Schmidt, A. Senger, M. Hauth, C. Freier, V. Schkolnik, and A. Peters. A mobile high-precision absolute gravimeter based on atom interferometry. *Gyroscopy and Navigation*, 2(3):170–177, jul 2011.
- [25] D. M. Harber, J. M. Obrecht, J. M. McGuirk, and E. A. Cornell. Measurement of the casimir-polder force through center-of-mass oscillations of a bose-einstein condensate. *Physical Review A*, 72(3), sep 2005.
- [26] Jonathan C. Knight. Photonic crystal fibres. *Nature*, 424(6950):847–851, aug 2003.
- [27] F. Benabid and P.J. Roberts. Linear and nonlinear optical properties of hollow core photonic crystal fiber. *Journal of Modern Optics*, 58(2):87–124, jan 2011.
- [28] Shoichi Okaba, Tetsushi Takano, Fetah Benabid, Tom Bradley, Luca Vincetti, Zakhar Maizelis, Valery Yampol'skii, Franco Nori, and Hidetoshi Katori. Lamb-dicke spectroscopy of atoms in a hollow-core photonic crystal fibre. *Nature Communications*, 5(1), jun 2014.
- [29] M. Bajcsy, S. Hofferberth, V. Balic, T. Peyronel, M. Hafezi, A. S. Zibrov, V. Vuletic, and M. D. Lukin. Efficient all-optical switching using slow light within a hollow fiber. *Physical Review Letters*, 102(20), may 2009.
- [30] Frank Blatt, Lachezar S. Simeonov, Thomas Halfmann, and Thorsten Peters. Stationary light pulses and narrowband light storage in a laser-cooled ensemble loaded into a hollow-core fiber. *Physical Review A*, 94(4), oct 2016.

-
- [31] M. R. Sprague, P. S. Michelberger, T. F. M. Champion, D. G. England, J. Nunn, X.-M. Jin, W. S. Kolthammer, A. Abdolvand, P. St. J. Russell, and I. A. Walmsley. Broadband single-photon-level memory in a hollow-core photonic crystal fibre. *Nature Photonics*, 8(4):287–291, mar 2014.
- [32] Vivek Venkataraman, Kasturi Saha, and Alexander L. Gaeta. Phase modulation at the few-photon level for weak-nonlinearity-based quantum computing. *Nature Photonics*, 7(2):138–141, dec 2012.
- [33] G. Epple, K. S. Kleinbach, T. G. Euser, N. Y. Joly, T. Pfau, P. St. J. Russell, and R. Löw. Rydberg atoms in hollow-core photonic crystal fibres. *Nature Communications*, 5(1), jun 2014.
- [34] C. L. Degen, F. Reinhard, and P. Cappellaro. Quantum sensing. *Reviews of Modern Physics*, 89(3), jul 2017.
- [35] D. Kielpinski, C. Monroe, and D. J. Wineland. Architecture for a large-scale ion-trap quantum computer. *Nature*, 417(6890):709–711, jun 2002.
- [36] T. Calarco, U. Dorner, P. S. Julienne, C. J. Williams, and P. Zoller. Quantum computations with atoms in optical lattices: Marker qubits and molecular interactions. *Physical Review A*, 70(1), jul 2004.
- [37] Y. O. Dudin, L. Li, and A. Kuzmich. Light storage on the time scale of a minute. *Physical Review A*, 87(3), mar 2013.
- [38] T. Kovachy, P. Asenbaum, C. Overstreet, C. A. Donnelly, S. M. Dickerson, A. Sugarbaker, J. M. Hogan, and M. A. Kasevich. Quantum superposition at the half-metre scale. *Nature*, 528(7583):530–533, dec 2015.
- [39] D. D. Yavuz, P. B. Kulatunga, E. Urban, T. A. Johnson, N. Proite, T. Henage, T. G. Walker, and M. Saffman. Fast ground state manipulation of neutral atoms in microscopic optical traps. *Physical Review Letters*, 96(6), feb 2006.
- [40] L. Isenhower, E. Urban, X. L. Zhang, A. T. Gill, T. Henage, T. A. Johnson, T. G. Walker, and M. Saffman. Demonstration of a neutral atom controlled-NOT quantum gate. *Physical Review Letters*, 104(1), jan 2010.
- [41] A. G. Radnaev, Y. O. Dudin, R. Zhao, H. H. Jen, S. D. Jenkins, A. Kuzmich, and T. A. B. Kennedy. A quantum memory with telecom-wavelength conversion. *Nature Physics*, 6(11):894–899, sep 2010.

- [42] Steven Chu. Nobel lecture: The manipulation of neutral particles. *Reviews of Modern Physics*, 70(3):685–706, jul 1998.
- [43] C Ryu and M G Boshier. Integrated coherent matter wave circuits. *New Journal of Physics*, 17(9):092002, sep 2015.
- [44] S. Kuhr, W. Alt, D. Schrader, I. Dotsenko, Y. Miroshnychenko, W. Rosenfeld, M. Khudaverdyan, V. Gomer, A. Rauschenbeutel, and D. Meschede. Coherence properties and quantum state transportation in an optical conveyor belt. *Physical Review Letters*, 91(21), nov 2003.
- [45] Pablo Solano, Jeffrey A. Grover, Jonathan E. Hoffman, Sylvain Ravets, Fredrik K. Fatemi, Luis A. Orozco, and Steven L. Rolston. Optical nanofibers. In *Advances In Atomic, Molecular, and Optical Physics*, pages 439–505. Elsevier, 2017.
- [46] Christos Markos, John C. Travers, Amir Abdolvand, Benjamin J. Eggleton, and Ole Bang. Hybrid photonic-crystal fiber. *Reviews of Modern Physics*, 89(4), nov 2017.
- [47] E. Vetsch, D. Reitz, G. Sagué, R. Schmidt, S. T. Dawkins, and A. Rauschenbeutel. Optical interface created by laser-cooled atoms trapped in the evanescent field surrounding an optical nanofiber. *Physical Review Letters*, 104(20), may 2010.
- [48] A. Goban, K. S. Choi, D. J. Alton, D. Ding, C. Lacroûte, M. Pototschnig, T. Thiele, N. P. Stern, and H. J. Kimble. Demonstration of a state-insensitive, compensated nanofiber trap. *Physical Review Letters*, 109(3), jul 2012.
- [49] D. Reitz, C. Sayrin, R. Mitsch, P. Schneeweiss, and A. Rauschenbeutel. Coherence properties of nanofiber-trapped cesium atoms. *Physical Review Letters*, 110(24), jun 2013.
- [50] Ashby P. Hilton, Christopher Perrella, Fetah Benabid, Ben M. Sparkes, Andre N. Luiten, and Philip S. Light. Exploring the limits of ultra-strong light-atom interaction in waveguide-trapped cold-atoms.
- [51] Maria Langbecker, Mohammad Noaman, Niels Kjærgaard, Fetah Benabid, and Patrick Windpassinger. Rydberg excitation of cold atoms inside a hollow-core fiber. *Physical Review A*, 96(4), oct 2017.
- [52] Mingjie Xin, Wui Seng Leong, Zilong Chen, and Shau-Yu Lan. An atom interferometer inside a hollow-core photonic crystal fiber. *Science Advances*, 4(1):e1701723, jan 2018.
- [53] C.S. Adams and E. Riis. Laser cooling and trapping of neutral atoms. *Progress in Quantum Electronics*, 21(1):1–79, jan 1997.

-
- [54] Metcalf Harold J. and Straten Peter. *Laser Cooling and Trapping of Neutral Atoms*. American Cancer Society, 2007.
- [55] Rudolf Grimm, Matthias Weidemller, and Yurii B. Ovchinnikov. Optical dipole traps for neutral atoms. In *Advances In Atomic, Molecular, and Optical Physics*, pages 95–170. Elsevier, 2000.
- [56] T.W. Hänsch and A.L. Schawlow. Cooling of gases by laser radiation. *Optics Communications*, 13(1):68–69, jan 1975.
- [57] Steven Chu, L. Hollberg, J. E. Bjorkholm, Alex Cable, and A. Ashkin. Three-dimensional viscous confinement and cooling of atoms by resonance radiation pressure. *Physical Review Letters*, 55(1):48–51, jul 1985.
- [58] H. J. Metcalf and P. van der Straten. Laser cooling and trapping of atoms. *Journal of the Optical Society of America B*, 20(5):887, may 2003.
- [59] A. Ashkin and J. P. Gordon. Stability of radiation-pressure particle traps: an optical earnshaw theorem. *Optics Letters*, 8(10):511, oct 1983.
- [60] Paul D. Lett, Richard N. Watts, Christoph I. Westbrook, William D. Phillips, Phillip L. Gould, and Harold J. Metcalf. Observation of atoms laser cooled below the doppler limit. *Physical Review Letters*, 61(2):169–172, jul 1988.
- [61] J. Dalibard and C. Cohen-Tannoudji. Laser cooling below the doppler limit by polarization gradients: simple theoretical models. *Journal of the Optical Society of America B*, 6(11):2023, nov 1989.
- [62] P. J. Ungar, D. S. Weiss, E. Riis, and Steven Chu. Optical molasses and multilevel atoms: theory. *Journal of the Optical Society of America B*, 6(11):2058, nov 1989.
- [63] William D. Phillips. Nobel lecture: Laser cooling and trapping of neutral atoms. *Reviews of Modern Physics*, 70(3):721–741, jul 1998.
- [64] Tomasz M. Brzozowski, Maria Maczynska, Michal Zawada, Jerzy Zachorowski, and Wojciech Gawlik. Time-of-flight measurement of the temperature of cold atoms for short trap-probe beam distances. *Journal of Optics B: Quantum and Semiclassical Optics*, 4(1):62, 2002.
- [65] O. Carnal and J. Mlynek. Young’s double-slit experiment with atoms: A simple atom interferometer. *Physical Review Letters*, 66(21):2689–2692, may 1991.

-
- [66] David W. Keith, Christopher R. Ekstrom, Quentin A. Turchette, and David E. Pritchard. An interferometer for atoms. *Physical Review Letters*, 66(21):2693–2696, may 1991.
- [67] Norman F. Ramsey. A molecular beam resonance method with separated oscillating fields. *Physical Review*, 78(6):695–699, jun 1950.
- [68] Norman F. Ramsey. *Molecular Beams*. Oxford University Press, London, 1956.
- [69] M. Kasevich and S. Chu. Measurement of the gravitational acceleration of an atom with a light-pulse atom interferometer. *Applied Physics B Photophysics and Laser Chemistry*, 54(5):321–332, may 1992.
- [70] Kathryn Moler, David S. Weiss, Mark Kasevich, and Steven Chu. Theoretical analysis of velocity-selective raman transitions. *Physical Review A*, 45(1):342–348, jan 1992.
- [71] A. Peters, K. Y. Chung, and S. Chu. High-precision gravity measurements using atom interferometry. *Metrologia*, 38(1):25, 2001.
- [72] K. Bongs, R. Launay, and M.A. Kasevich. High-order inertial phase shifts for time-domain atom interferometers. *Applied Physics B*, 84(4):599–602, aug 2006.
- [73] B. Young, M. Kasevich, and S. Chu. *Precision Atom Interferometry with Light Pulses*. Academic Press, San Diego, CA, 1997.
- [74] F. Benabid. Stimulated raman scattering in hydrogen-filled hollow-core photonic crystal fiber. *Science*, 298(5592):399–402, oct 2002.
- [75] F. Couny, F. Benabid, P. J. Roberts, P. S. Light, and M. G. Raymer. Generation and photonic guidance of multi-octave optical-frequency combs. *Science*, 318(5853):1118–1121, nov 2007.
- [76] B. Debord, M. Alharbi, T. Bradley, C. Fourcade-Dutin, Y.Y. Wang, L. Vincetti, F. Gérôme, and F. Benabid. Hypocycloid-shaped hollow-core photonic crystal fiber part i: Arc curvature effect on confinement loss. *Optics Express*, 21(23):28597, nov 2013.
- [77] <http://www.glophotronics.fr/index.php/hollow-core-photonic-crystal-fibers.html>.
- [78] Michal Bajcsy. *Novel Systems and Techniques for Nonlinear Optics at Low Light Levels*. PhD thesis, Harvard University, 2009.
- [79] Carl E. Wieman and Leo Hollberg. Using diode lasers for atomic physics. *Review of Scientific Instruments*, 62(1):1–20, jan 1991.

-
- [80] Xuejian Wu, Fei Zi, Jordan Dudley, Ryan J. Bilotta, Philip Canoza, and Holger MÃ(eller. Multiaxis atom interferometry with a single-diode laser and a pyramidal magneto-optical trap. *Optica*, 4(12):1545, dec 2017.
- [81] Zilong Chen, Justin G. Bohnet, Joshua M. Weiner, and James K. Thompson. A low phase noise microwave source for atomic spin squeezing experiments in 87rb. *Review of Scientific Instruments*, 83(4):044701, apr 2012.
- [82] Norman F. Ramsey. Experiments with separated oscillatory fields and hydrogen masers. *Reviews of Modern Physics*, 62(3):541–552, jul 1990.
- [83] P. Hamilton, M. Jaffe, P. Haslinger, Q. Simmons, H. Muller, and J. Khoury. Atom-interferometry constraints on dark energy. *Science*, 349(6250):849–851, aug 2015.
- [84] Holger Mller, Achim Peters, and Steven Chu. A precision measurement of the gravitational redshift by the interference of matter waves. *Nature*, 463(7283):926–929, feb 2010.
- [85] D. Schlippert, J. Hartwig, H. Albers, L. L. Richardson, C. Schubert, A. Roura, W. P. Schleich, W. Ertmer, and E. M. Rasel. Quantum test of the universality of free fall. *Physical Review Letters*, 112(20), may 2014.
- [86] Lin Zhou, Shitong Long, Biao Tang, Xi Chen, Fen Gao, Wencui Peng, Weitao Duan, Jiaqi Zhong, Zongyuan Xiong, Jin Wang, Yuanzhong Zhang, and Mingsheng Zhan. Test of equivalence principle at10-8level by a dual-species double-diffraction raman atom interferometer. *Physical Review Letters*, 115(1), jul 2015.
- [87] I. Dutta, D. Savoie, B. Fang, B. Venon, C. L. Garrido Alzar, R. Geiger, and A. Landragin. Continuous cold-atom inertial sensor with1 nrad/secRotation stability. *Physical Review Letters*, 116(18), may 2016.
- [88] M. de Angelis, A. Bertoldi, L. Cacciapuoti, A. Giorgini, G. Lamporesi, M. Prevedelli, G. Saccorotti, F. Sorrentino, and G. M. Tino. Precision gravimetry with atomic sensors. *Measurement Science and Technology*, 20(2):022001, 2009.
- [89] T. A. Savard, K. M. O’Hara, and J. E. Thomas. Laser-noise-induced heating in far-off resonance optical traps. *Physical Review A*, 56(2):R1095–R1098, aug 1997.
- [90] I. Dotsenko, W. Alt, S. Kuhr, D. Schrader, M. Mller, Y. Miroshnychenko, V. Gomer, A. Rauschenbeutel, and D. Meschede. Application of electro-optically generated light fields for raman spectroscopy of trapped cesium atoms. *Applied Physics B*, 78(6):711–717, apr 2004.

-
- [91] L. P. Parazzoli, A. M. Hankin, and G. W. Biedermann. Observation of free-space single-atom matter wave interference. *Physical Review Letters*, 109(23), dec 2012.
- [92] S. Marksteiner, C. M. Savage, P. Zoller, and S. L. Rolston. Coherent atomic waveguides from hollow optical fibers: Quantized atomic motion. *Physical Review A*, 50(3):2680–2690, sep 1994.
- [93] Brynle Barrett, Rémy Geiger, Indranil Dutta, Matthieu Meunier, Benjamin Canuel, Alexandre Gauguier, Philippe Bouyer, and Arnaud Landragin. The sagnac effect: 20 years of development in matter-wave interferometry. *Comptes Rendus Physique*, 15(10):875–883, dec 2014.
- [94] H. J. Kimble. The quantum internet. *Nature*, 453(7198):1023–1030, jun 2008.
- [95] Klemens Hammerer, Anders S. Sørensen, and Eugene S. Polzik. Quantum interface between light and atomic ensembles. *Reviews of Modern Physics*, 82(2):1041–1093, apr 2010.
- [96] Ekaterina Ilinova, James F. Babb, and Andrei Derevianko. Feasibility of an optical fiber clock. *Physical Review A*, 96(3), sep 2017.
- [97] Maria Langbecker, Mohammad Noaman, Niels Kjærgaard, Fetah Benabid, and Patrick Windpassinger. Rydberg excitation of cold atoms inside a hollow-core fiber. *Physical Review A*, 96(4), oct 2017.
- [98] L. Allen and J. H. Eberly. Optical resonance and two level atoms. *Physik in unserer Zeit*, 7(1):31–32, 1976.
- [99] Jiaheng Yang, Xiaodong He, Ruijun Guo, Peng Xu, Kunpeng Wang, Cheng Sheng, Min Liu, Jin Wang, Andrei Derevianko, and Mingsheng Zhan. Coherence preservation of a single neutral atom qubit transferred between magic-intensity optical traps. *Physical Review Letters*, 117(12), sep 2016.
- [100] A. W. Carr and M. Saffman. Doubly magic optical trapping for cs atom hyperfine clock transitions. *Physical Review Letters*, 117(15), oct 2016.
- [101] Chang Huang, Pei-Chen Kuan, and Shau-Yu Lan. Laser cooling of ^{85}rb atoms to the recoil-temperature limit. *Physical Review A*, 97(2), feb 2018.
- [102] G. Rosi, F. Sorrentino, L. Cacciapuoti, M. Prevedelli, and G. M. Tino. Precision measurement of the newtonian gravitational constant using cold atoms. *Nature*, 510(7506):518–521, jun 2014.

-
- [103] Michael A. Hohensee, Brian Estey, Paul Hamilton, Anton Zeilinger, and Holger Müller. Force-free gravitational redshift: Proposed gravitational aharonov-bohm experiment. *Physical Review Letters*, 108(23), jun 2012.
 - [104] Onur Hosten, Nils J. Engelsen, Rajiv Krishnakumar, and Mark A. Kasevich. Measurement noise 100 times lower than the quantum-projection limit using entangled atoms. *Nature*, 529(7587):505–508, jan 2016.
 - [105] Kevin C. Cox, Graham P. Greve, Joshua M. Weiner, and James K. Thompson. Deterministic squeezed states with collective measurements and feedback. *Physical Review Letters*, 116(9), mar 2016.
 - [106] E. Knill, R. Laflamme, and G. J. Milburn. A scheme for efficient quantum computation with linear optics. *Nature*, 409(6816):46–52, jan 2001.
 - [107] L.-M. Duan, M. D. Lukin, J. I. Cirac, and P. Zoller. Long-distance quantum communication with atomic ensembles and linear optics. *Nature*, 414(6862):413–418, nov 2001.
 - [108] Morgan P. Hedges, Jevon J. Longdell, Yongmin Li, and Matthew J. Sellars. Efficient quantum memory for light. *Nature*, 465(7301):1052–1056, jun 2010.
 - [109] R. Zhao, Y. O. Dudin, S. D. Jenkins, C. J. Campbell, D. N. Matsukevich, T. A. B. Kennedy, and A. Kuzmich. Long-lived quantum memory. *Nature Physics*, 5(2):100–104, dec 2008.