

**An On-Chip Atom Interferometer Using a
Bose-Einstein Condensate**

by

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An On-Chip Atom Interferometer Using a Bose-Einstein Condensate

Thesis directed by Prof. Eric A. Cornell

This work reports the first demonstration of an on-chip atom Michelson interferometer employing a Bose-Einstein condensate (BEC). An intra-waveguide optical standing wave serves to split, reflect, and recombine the BEC while the BEC is confined and propagates in a magnetic waveguide. The combined optical beamsplitter and magnetic waveguide employed here have allowed us to study the coherence properties of matter waves confined in a microstructure.

To prove that the interferometer is successfully working, a differential phase shift is introduced between the two arms of the interferometer with a magnetic field gradient and alternatively with an initial condensate velocity in a trap with a longitudinal frequency of 5 Hz. We observe interference when the round-trip propagation time is relatively short, i.e., less than about 10 ms, and the maximum separation of the split wave packets is about 120 μm .

This thesis also introduces an alternative way of making a BEC by surface-induced evaporative cooling. The surface-induced cooling provides a promising technique, which can be potentially used for a continuous coherent source on the chip.

Dedication

To my family for all of their love and support.

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My graduate study at CU and JILA was doubtless the most rewarding experience because of the wonderful people who I had the privilege to work with. I am grateful to have Eric Cornell as my advisor. His enthusiasm for research and life has been a great inspiration and his physics insight has always amazed me and keep me motivated. I would also like to thank my second advisor, Dana Anderson, who never hesitated to give me help and offer me guidance when I knocked his door. He always encouraged me and taught me how to become a better person and scientist through my graduate career. I also appreciate Carl Wieman and Deborah Jin for their useful comments and advises.

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Chapter 1

Introduction

Over the past several years, the manipulation of atomic matter waves has become a dominant theme for atomic physics experiments. The creation of the first Bose-Einstein Condensate (BEC) in 1995 by E. Cornell and C. Wieman at JILA [1] and W. Ketterle at MIT [2] has opened up a whole new field of research, namely the study of coherent matter waves. In the same way as coherent light sources in the form of lasers have revolutionized modern optics, the manipulation of coherent matter waves is creating a deep impact on atomic physics. In fact, the analogy between light waves and matter waves is very fundamental and leads to the idea of atom interferometry, i.e., using atomic matter waves for measurement of rotation in a similar way as lasers are used in laser gyroscopes.

The measurement of rotation rates using a laser gyroscope is based on the Sagnac effect, in which a phase difference between two interferometer arms arises due to a rotation, whose rotating axis is perpendicular to the interferometer plane. Macek and Davis firstly recognized in 1963 [3] that photons propagating in different arms see different path lengths when the rotation happens, thus carry a phase difference, which is given by

$$\Delta\phi = \vec{A} \cdot \vec{\Omega}_{rot} \frac{E_r}{\hbar c^2}, \quad (1.1)$$

where $\vec{A}$ is the enclosed area of the interferometer, $\vec{\Omega}_{rot}$ is the rotation rate, $E_r = \hbar\omega$ is the energy of the photons, and c is the speed of the light. By detecting this phase

difference, the rotating motion can thus be measured. Today, optical gyroscopes, such as ring laser gyros or fiber gyros, are widely used as precise inertial sensors for the navigation.

Atom interferometry offers an exquisite precision measurement capability with far greater sensitivity than its photon-based counterpart. For example, the fundamental limit on the signal-to-noise ratio of an atom Sagnac gyroscope is a factor of 10^{11} greater than the optical one, given comparable enclosed areas and particle flux. This amazing improvement is due to the high relativistic energy of an atom, mc^2 , as compared to the photon energy, $E_r = \hbar\omega$. Besides, atoms can be sensitive to electric and magnetic fields. Because of this, atom interferometry may be suited to a substantially larger number of sensor applications despite being sensitive to more detrimental noise sources. Atom interferometry experiments have revealed promising and sometimes stunning measurement capabilities [4, 5, 6, 7, 8].

Experimental research in the field of BEC has been dealing with two different aspects. First, experiments have investigated and are continuing to explore the fundamental properties of BECs. Second, new technological developments allow us to create and manipulate BECs in a more and more flexible way. One key element here is the development of “atom chips”.

Atom chips seek to implement atom-based devices on a small scale [9, 10, 11, 12, 13]. For example, one can incorporate conductors, magnetic elements, and optical components on a single substrate to produce fields that confine, control, and manipulate atoms. One can also incorporate atom detection and signal conversion on that same substrate. Moreover, because the source of the magnetic fields on a chip are close to the atoms, one can typically apply much larger forces (or use much less power to apply those forces). Most importantly, the atom chip approach offers substantial control over the geometry of an atom interferometer.

Most attempts to implement a coherent beamsplitter/recombiner, which is es-

essential for an interferometer, on a chip have used current-induced magnetic fields to split a condensate cloud. The magnetic field typically forms a double potential well that merges and then splits apart either in space, in time, or in both. However, atom-atom interactions due to high density in the waveguide can cause instabilities during the splitting or recombining process in such a double-well potential [14] and also reduce the phase coherence [15]. Various detrimental atom-surface and atom-wire interactions [16, 17, 18, 19, 20] have been reported. Because of those interactions, the coherence of the condensate is destroyed during the splitting process and attempts to demonstrate on-chip interference have been stymied. However, optical beamsplitters using standing light waves have offered promising potential for coherent manipulation of atoms [6, 7, 8]. The combined optical and magnetic forces have enabled a study of coherence properties of matter waves confined in a microstructure.

This thesis work brings together the two aspects of modern BEC research by exploring coherent processes of light-matter interaction and implementing them on an atom chip to demonstrate the first working on-chip atom interferometer. This work is done in collaboration with Prof. M. Prentiss at Harvard University.

In this thesis, I will show the first on-chip atom interferometer using a standing-light wave. The details of the atom chip and the apparatus are described in Chapter 2. In Chapter 3, we also introduce an alternative way of making a condensate by surface-induced evaporative cooling. The surface-induced cooling provides a promising technique, which can potentially be used to produce a continuous-wave (cw) Bose-Einstein Condensate (BEC) on a chip. Coherent splitting and reflecting of the condensates by diffraction from a standing light wave are discussed in Chapter 4. An innovative double-pulse beamsplitter, which offers a simple solution for coherently splitting condensates, is described in Chapter 5. In Chapter 6, we develop an on-chip Michelson interferometer to split, reflect, and recombine the condensate and read out interference signals. To prove that the interferometer is working successfully, a differential phase shift is created

by a magnetic gradient and alternatively by an initial velocity in a trap with a longitudinal frequency of 5 Hz. This robust interferometer technique could potentially be used for an atom gyroscope.

In summary, this thesis demonstrates the first realization of a guided atom interferometer on an atom chip, which has been a long-standing goal of research in this field. The experiments described in this thesis are successful “proof-of-principle” experiments, providing a way for future applications of BEC in the field of ultra-high precision measurement.

Chapter 2

Experimental setup

Our experiment employs a modular system that consists of three adjoint vacuum chambers and an atom chip on which a BEC and an interferometer are made [20]. The multiple-chamber design is based on the fact that the vacuum requirements are different at different cooling stages. In each chamber, the vacuum is sustained by its own pumping system, which gives us the freedom to control the pressure separately. The overview of this modular system and the cooling process associated with each chamber is described in Section 2.1. I will discuss the atom chip that is used to create a condensate and an interferometer in particular in Section 2.2. The atom chip consists of lithographically-patterned wires and an integrated prism pair. The wires are used to produce the magnetic fields for atom trapping/guiding, and the prism pair is used to produce an optical standing wave for a beamsplitter. At the end of this chapter, Section 2.3, I will describe the optical setup used to project a laser beam close to the chip surface to form the optical standing wave.

2.1 Overview of the apparatus for producing a Bose-Einstein Condensate

The apparatus consists of the pyramid MOT chamber, pre-evaporation chamber and application chamber, as shown in Figure 2.1 [20]. The pyramid MOT chamber includes a pyramid mirror and a pair of coils to provide a magneto-optical trap (MOT)

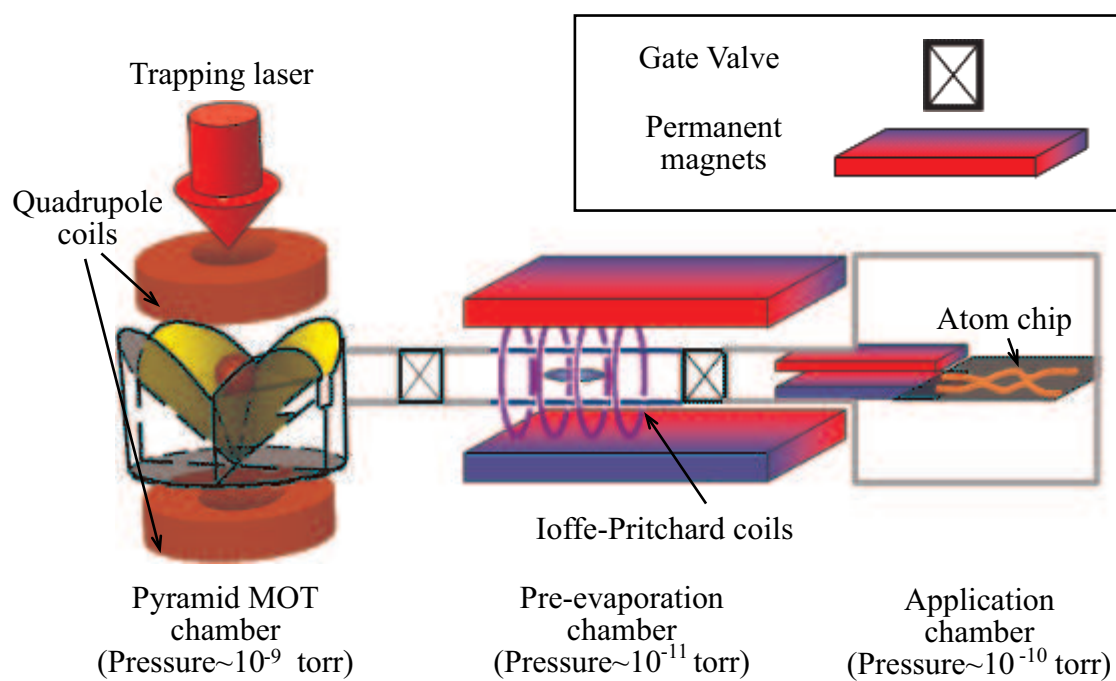


Figure 2.1: Schematic of the apparatus. The apparatus consists of three chamber separated by pumping stages to support pressure requirements for different cooling stages.

for laser cooling and trapping. A 200 mW single laser beam with a diameter of approximately 10 cm is projected onto a pyramid mirror, which sits at the bottom of the MOT chamber. Reflections from the pyramid mirror construct the six-beam configuration required for laser cooling. Atoms that are in the weak-field seeking states would be trapped in the minimum of a magnetic trap so we use the coil pair in an anti-Helmholtz configuration generates a quadrupole trap to confine the cold atoms. The coils sit on a track and can be mechanically moved down to the second chamber to transfer atoms from the pyramid MOT chamber to the pre-evaporation chamber. A pair of permanent magnets and four coils in an Ioffe-Pritchard (IP) configuration in the pre-evaporation chamber generate magnetic fields to form a three-dimensional harmonic trap, a so called “IP trap”, for evaporative cooling. In the application chamber, an atom chip that contains lithographically patterned wires is installed. The wires generate magnetic fields that are essential for a microtrap and a magnetic waveguide. A pair of prisms is integrated on the chip to create a standing light wave needed for a coherent beamsplitter.

Because different stages of cooling require different vacuum quality, we built this modular vacuum system to allow us to modify the pressure in each chamber individually. Every chamber is connected to a separate ion pump and the conductances between the chambers are negligible compared to the pumping speed, so that the pressure in each chamber can be controlled separately. The pre-evaporation chamber and the application chamber also have Titanium sublimation pumps (TSPs) installed for better pumping speeds so that the atom cloud has a longer lifetime in those two chambers to support evaporative cooling. Each two adjoint chambers are connected by gate valves, which allows us to isolate one of the chambers and modify its pressure separately. For example, by closing the gate valve in between the pre-evaporation chamber and application chamber, we can vent, open the application chamber to replace a chip, and pump down again after closing it without affecting the ultrahigh vacuum of the pre-evaporation chamber. This modular system allows us to replace a atom chip and prepare the vacuum typically

in 4–5 days. Additional details of our apparatus and BEC production can be found in Ref. [20].

To make a condensate, Rb atoms are first collected and trapped from a vapor pressure of 10^{-9} torr in the pyramid MOT chamber. After the atoms are cooled down to around 500 μK , they are loaded into the magnetic quadrupole trap. Atoms are then mechanically moved by moving the quadrupole coils to the pre-evaporation chamber, where they are transferred into the IP trap. Rf-forced evaporative cooling is applied to cool the atoms down to a temperature just above the critical temperature for forming a BEC. To promote the efficiency of the evaporative cooling, the pre-evaporation chamber must have a low pressure, i.e., in the range of 10^{-11} torr. After evaporative cooling, the cold atoms slide down a potential slope produced by a magnetic gradient, and move toward the atom chip, which is located in the application chamber. We capture and load atoms into a microtrap located on the chip surface. Another stage of rf-forced evaporative cooling is applied to form a BEC on the chip.

2.2 Atom chip with integrated prisms

Our atom chip, the dimensions of which are 5 cm by 2 cm, contains lithographically patterned copper wires and a pair of prism-shaped mirrors integrated on an aluminum nitride substrate, as shown in Figure 2.2. The wires generate the magnetic fields needed to create the microtrap and magnetic waveguide. The tapered wires at the entrance side help to capture and load atoms onto the chip. The straight wire and a transverse bias field (B_t) perpendicular to it provide a magnetic waveguide. The L-shaped wire in combination with a longitudinal bias field (B_l) creates the longitudinal confinement to form a microtrap.

The transverse bias field cancels out the magnetic field created by the straight wire, called “guiding wire”, to form a magnetic potential minimum above the wire, therefore, a two dimensional confinement is created along the wire, as shown in Figure

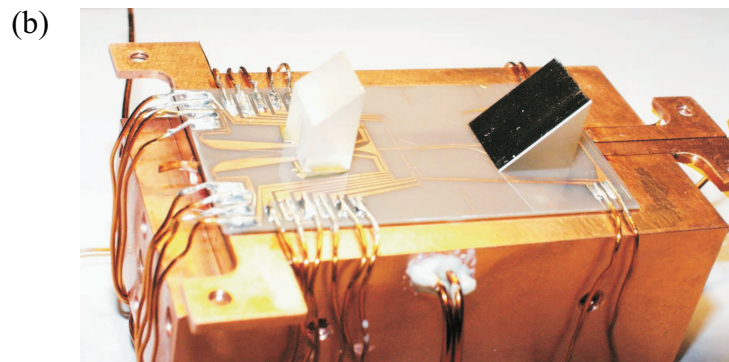
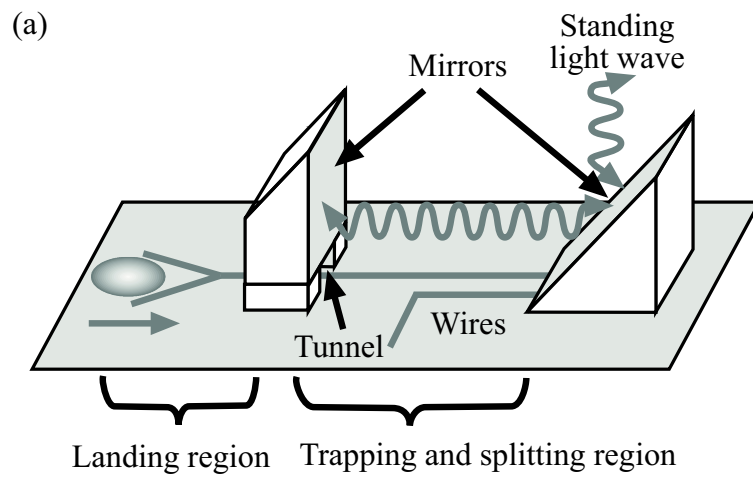


Figure 2.2: (a) Schematic of the atom chip (not to scale). The prism-shaped mirrors are integrated with microfabricated wires on an aluminum nitride substrate. The dimensions of the whole chip are 5 cm by 2 cm. (b) A photo image of the atom chip.

2.3. The minimum of the field extends along the guiding wire and atoms can be trapped in two dimensions but are free to move along the wire to form a magnetic waveguide. The distance (D) of the waveguide from the center of the wire is given by

$$D = \frac{\mu_0}{2\pi} \frac{I}{B_t}, \quad (2.1)$$

where I is the current in the guiding wire and μ_0 is the magnetic permeability. A three dimensional microtrap can be created by adding a second “trapping wire”, perpendicular to the first wire and also a longitudinal bias field, as shown in Figure 2.4. In our experiment, we use a “L”-shaped wire that has a perpendicular component to the guiding wire to create the longitudinal confinement.

The inward-facing surfaces of the two prisms are mirror-coated and aligned perpendicular to the magnetic waveguide to produce optical standing waves. When the atoms are loaded onto the chip, they are typically $115 \mu\text{m}$ away from the surface so a $180\text{-}\mu\text{m}$ -high tunnel is created underneath the prism that is located at the entrance side of the chip to allow for loading the atoms in the microtrap. The cloud is delivered through the tunnel and captured in the chip’s trapping region, which is located between the prisms and sits $115 \mu\text{m}$ away from the surface of the chip. The atoms are then brought farther away from the surface to the center of the standing light wave, which is $250 \mu\text{m}$ above the surface to avoid the tunnel. Next, rf-forced evaporative cooling is used to form a condensate. The condensate is released into the magnetic waveguide. We turn on an optical standing wave close to the surface of the chip that splits, reflects, and recombines the atom clouds to form an atom interferometer.

The optical standing wave is generated by two 45-degree-cut prisms on the surface of the chip. The prisms are commercial BK7 right-angle prisms. The dimensions of the prisms are $1 \text{ cm} \times 1 \text{ cm} \times 1 \text{ cm}$ and the angle tolerance is less than 3 arc min. The prisms are glued on to the the chip surface with vacuum-compatible epoxy. The inward surface of each prism is mirror-coated. A tightly focused beam is incident on and reflected by

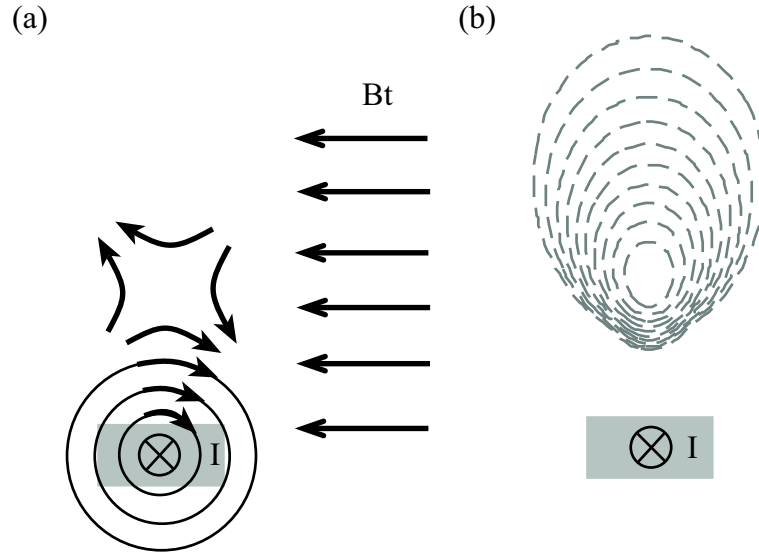


Figure 2.3: A magnetic waveguide. (a) The solid lines represent the magnetic fields created by the transverse bias field and the wire. The transverse bias field (B_t) cancels out the magnetic field created by the straight wire at one point. A magnetic potential minimum is formed above the wire. (b) The dashed lines are the contours of the magnetic field. The inner contours represent lower magnetic fields.

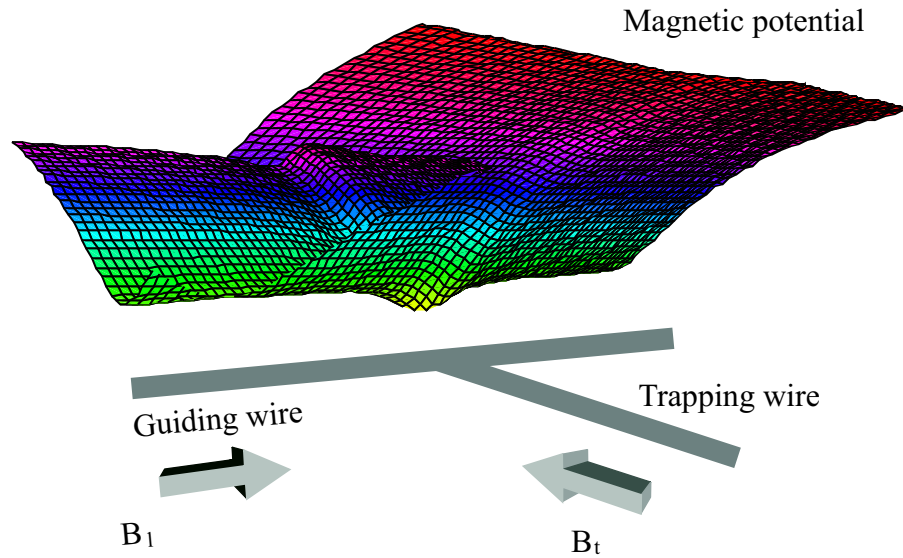


Figure 2.4: The magnetic potential of the microtrap. The three dimensional microtrap is created by two perpendicular wires and two perpendicular bias fields.

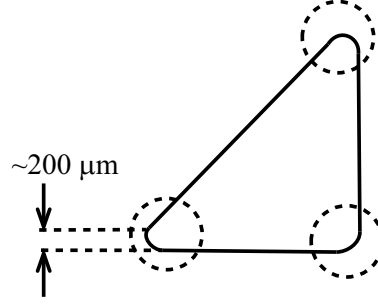


Figure 2.5: A bevelled prism (side view) (not to scale). The prism edges (marked by dotted circles) are bevelled to prevent them from chipping.

the mirror that is oriented at 45 degrees with respect to the chip surface. The reflected beam is parallel to the surface of the chip and the direction of the waveguide. The beam is retroreflected by the second mirror that is perpendicular to the waveguide and the surface of the chip. The technical details of the prisms are described in the following sections.

2.2.1 Preparation and installation of the prisms on the chip

We prepare the prisms by lapping, coating, and cleaning them before installing them onto the chip. The prisms first need to be lapped down by 100–200 μm . Lapping is necessary because commercial prisms are usually bevelled on the edges (see Figure 2.5) to prevent damages caused by stress during handling. Those bevelled edges cause nonuniform reflections of our optical beam because the optical beam is reflected very close to the chip surface, i.e., the edge of the prisms.

The inward surfaces of the two prisms (see Figure 2.2) are mirror-coated. The polarization and intensity of the reflected beam are determined by the reflectivity of the coating and any phase shift introduced by it. To have the best contrast ratio of the standing light wave, the polarization and the intensity of the incident and reflected beams must be the same. Therefore, we choose silver as our mirror coating because it minimizes polarization distortion. Silver coating also has a high reflectivity, typically

greater than 96% at a wavelength of 780 nm. The first step in the coating process is to coat the surface with a silver layer of thickness typically 150–200 nm. A thin protective layer of dielectric material, silicon dioxide (SiO_2), is then overlayed to protect the metallic-mirror coating from abrasion. The SiO_2 is sufficiently thin that the phase shift introduced by it is negligible. We clean the prisms carefully with Acetone and Dichloromethane after lapping and coating to ensure vacuum compatibility. We then install the prisms onto the chip.

The first step in installing the two prisms is to apply a thin layer of vacuum-compatible epoxy (Epoxy Technology, product # 353ND) to the bottom surfaces of the prisms. Because a tunnel has to be created underneath the prism at the entrance of the atom chip to allow atoms going through, two small flat pieces of glass with a thickness of 180 μm and a surface area of approximately 2.5 mm $\times$ 5 mm serve as spacers between the prism and the chip. Second, as shown in Figure 2.6, we align the prism precisely to the marks on the chip using a microscope. The marks are made with copper (Cu) wires during the lithographic production of the chip. Because the marks are made at the same time as the guiding and trapping wires, the alignment of the marks with respect to the waveguide should be good within 1 μm . Finally, the prisms are held down on the chip with a customized mount after the alignment and the chip assembly is put into an oven and heated to 100° for an hour to cure the epoxy. Because prism alignment is the most crucial step in this process, we will discuss it in detail in the following section.

2.2.2 Alignment of the prisms

The alignment of the prisms is important because the propagation direction of the standing-wave light field with respect to the waveguide affects the transverse momentum of the atoms.

The wave vector of the standing-wave light field needs to be parallel to the waveguide to minimize radial excitations of the BEC from photon scattering. The tolerance

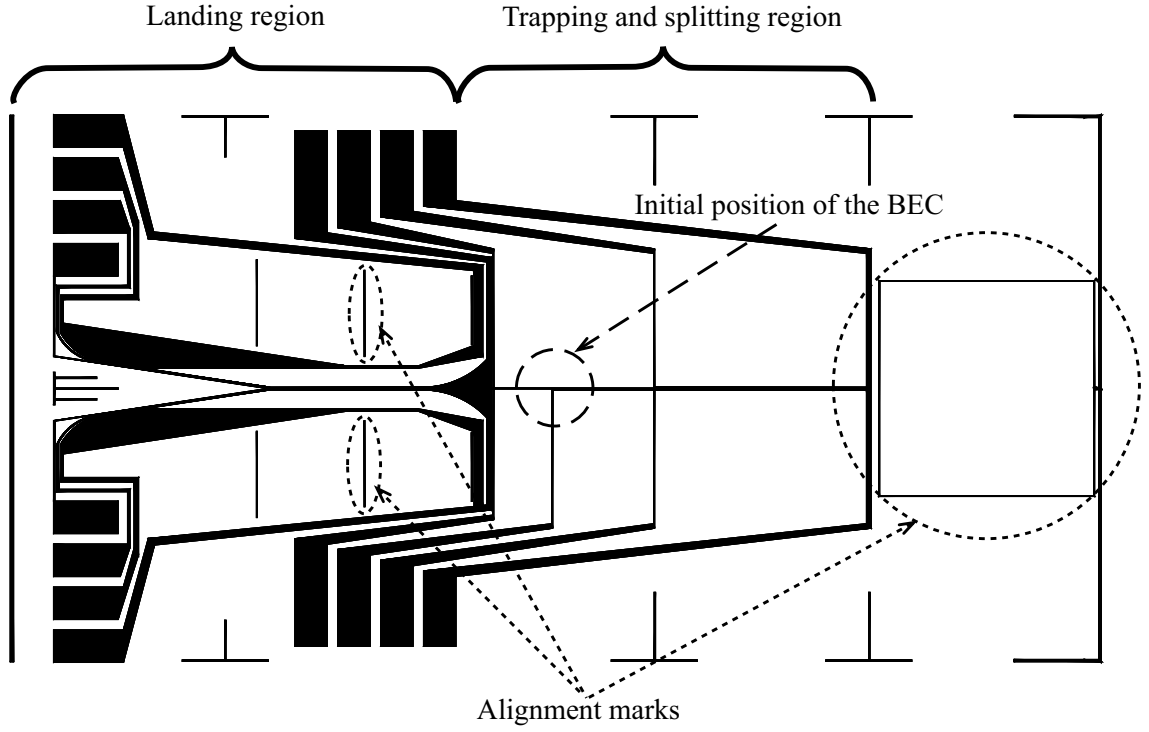


Figure 2.6: Alignment marks for the prism. All the marks (circled by the dotted lines) are $20\text{-}\mu\text{m}$ wide and $10\text{-}\mu\text{m}$ high. The dashed circle indicates the position where the Bose-Einstein Condensate (BEC) is formed.

of the alignment of the prisms with respect to the waveguide can be calculated by comparing the transverse momentum kick (δp) due to photon scattering to the transverse momentum (p_0) of atoms in the ground state. To avoid radial excitation, we need the condition:

$$\delta p \ll p_0. \quad (2.2)$$

We assume that the angle of the optical standing wave due to the misalignment with respect to the waveguide is $\delta\theta$. The transverse momentum kick from a photon-scattering process is

$$\delta p = n\hbar k_0 \sin \delta\theta, \quad (2.3)$$

where n is the number of scattered photons and k_0 is the wave vector of the photon. The root-mean-square (rms) of the momentum of non-interacting atoms in the ground state of a harmonic potential can be written as

$$p_0 = \sqrt{\hbar m_{Rb} \omega / 2}, \quad (2.4)$$

where m_{Rb} is the mass of the atoms and ω is the radial trap frequency. For a given p_0 in Equation 2.4, the angle tolerance of the misalignment $\delta\theta$ is smaller, when n is greater. The highest order of the diffraction in our experiment is a four-photon scattering process, which is used to reflect split clouds. Thus, for $n=4$, if we substitute δp and p_0 in Equation 2.2 with Equations 2.3 and 2.4, we have the angle tolerance of the misalignment with respect to the waveguide as

$$\delta\theta \ll \sin^{-1} \left(\frac{\sqrt{m_{Rb} \omega / 2 \hbar}}{4k_0} \right). \quad (2.5)$$

For a radial trap frequency of 100 Hz and a wavelength of 780 nm, the angle tolerance of the misalignment is estimated to be 1.2° from equation 2.5.

The alignment of the prisms can be done within the required accuracy. We achieve a position accuracy within the line width (δd) of the marks, i.e., $20 \mu\text{m}$, by aligning the edges of the prisms parallel to the marks under the microscope. The width (D) of the

prisms is 1 cm so we have a maximum for the tilted angle less than $\sin^{-1}(\delta d/D) = 0.11^\circ$ after the alignment. The alignment is also checked after the installation by shining a laser beam with a waist of about 2–3 mm onto the prisms. We observe a good overlap of the incident and reflected beams at long distance, e.g., 2–3 m, from the chip. We also make sure the prisms are aligned parallel to the mark within the same accuracy.

One other concern is the angle tilt introduced by the lapping process if a nonuniform layer is removed during the process. During lapping, a layer of maximum 200 μm thickness is taken off over a length of 1 cm. If the lapping process wasn't done carefully, one could get a tilt angle of $\tan^{-1}(\frac{200\mu\text{m}}{1\text{cm}}) = 1.27^\circ$ in the worst case. However, our lapping process can control the thickness with a precision of less than 100 μm so that the tilt angle should be within the tolerance.

2.3 Optical setup of a standing light wave close to the surface

The precisely aligned prisms sit on the atom chip in the vacuum chamber and the optical beam needs to be brought to the prism pair from outside the vacuum chamber so we design an optical system to shape the optical beam and project it onto the prisms (see Figure 2.7).

A frequency-stabilized external-cavity diode laser provides the light source for the standing wave. The frequency of the laser beam is locked at the ^{87}Rb D2 line, $|F = 2\rangle \longrightarrow |F = 3\rangle$ transition with a frequency instability of less than 1 MHz. Because atoms are in the $|F = 1\rangle$ state, this frequency is 6.8-GHz red-detuned from the atomic transition. The laser beam passes through an acousto-optical modulator and we adjust the acousto-optical modulator to have maximum power in the 1st order output. The other output orders are blocked by a pinhole. The acousto-optical modulator is used as a fast switch to create short pulses with pulse widths typically from a few μs to a few hundred μs . The center frequency of the acousto-optical modulator is 80 MHz so that the frequency detuning introduced by the acousto-optical modulator is negligible

compared to -6.8 GHz. The first order output of the is coupled into a single-mode polarization-maintaining fiber to preserve the polarization stability. The output of the fiber is shaped into a waist of $110\text{ }\mu\text{m}$ and we inject the beam into the vacuum chamber from the bottom window.

Note that a Gaussian beam has a divergence angle and therefore, the beam also has a non-zero transverse momentum spread due to its finite waist. For a beam waist of $110\text{ }\mu\text{m}$ and a wavelength of 780 nm , the divergence angle is 0.129° , which is well within the acceptable tolerance.

The output of the fiber is collimated by the first lens (L1) and passes through two wave plates: a half-wave plate and a quarter-wave plate. The two wave plates set the polarization of the beam to be linear and parallel to the chip surface. After the wave plates, the beam is partially reflected by a 8° wedge. This 5% reflection is collected by a photodetector, which is used to monitor the intensity of the laser beam. The beam is shaped by three lenses: L2, L3, and L4. The focal lengths of the three lenses are -25, 75.6, and 500 mm, respectively. Distances between the lenses are 50 and 15 mm. The first two lenses, L2 and L3, form a telescope that scales the beam from an original waist of approximately $200\text{ }\mu\text{m}$ to a waist of 1.13 mm . A concave lens is used instead of a convex lens to decrease the required distance of the telescope. The last lens, L4, focuses the size down to a waist of $110\text{ }\mu\text{m}$. The waist of the beam is positioned at the surface of the second prism corresponding to a Rayleigh length of 4.9 cm , where the Rayleigh range is the distance in which the beam size increases by a factor of $\sqrt{2}$ [21]. Because the Rayleigh range is much larger than the distance between the BEC and the second mirror, i.e., 2 mm , the beam can be considered collimated between the prisms, and the wave vector of the beam parallel to the waveguide. Mirror M2 sits on a translation stage that allows us to change beam height from the chip surface.

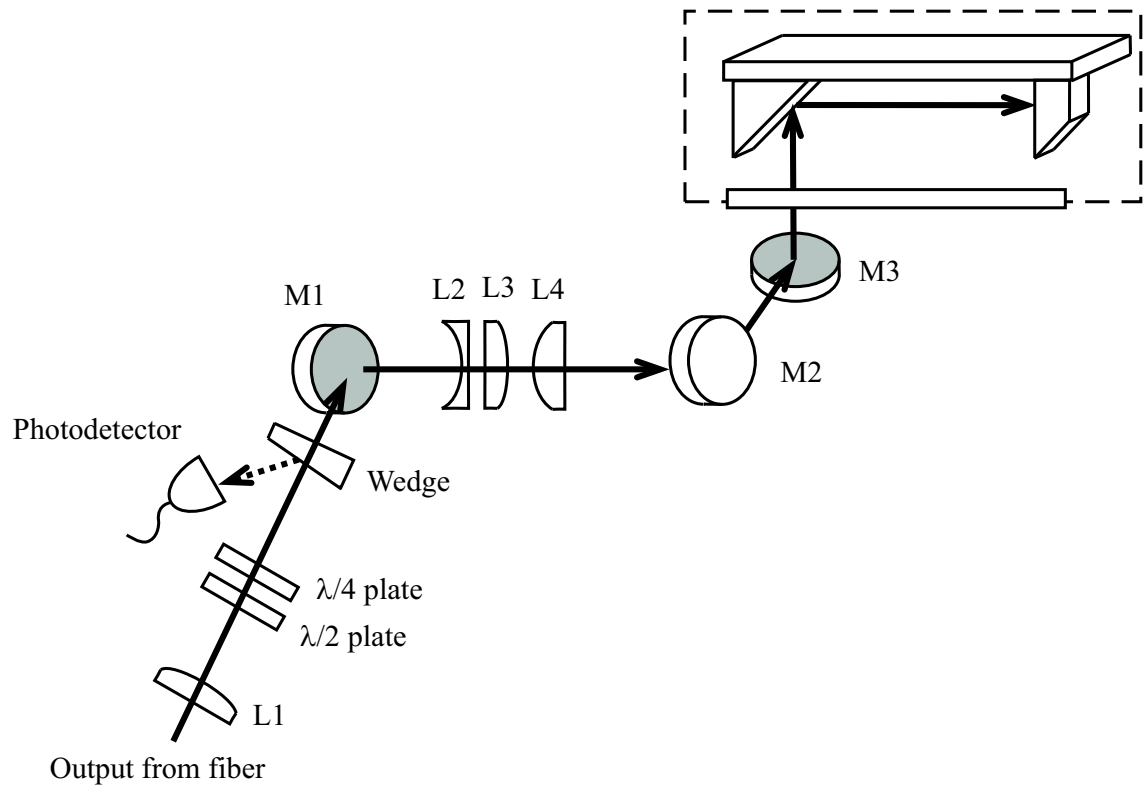


Figure 2.7: Schematic of the optical layout for a standing wave on a chip surface (not to scale). L1, L2, L3 and L4 are lenses. M1, M2, and M3 are metallic mirrors. The output beam from the fiber is collimated by L1. L2, L3, and L4 focus the beam to a waist of $110\ \mu\text{m}$. The waist of the beam is located at the surface of the second prism. Two wave plates create a linear polarization parallel to the chip surface. Approximately 5% of the output power from the fiber is reflected by a wedge surface and collected by the photodiode to monitor the power of the standing light wave.

Chapter 3

Surface-induced evaporative cooling

Conventional evaporative cooling often employs rf radiation. However, evaporative cooling induced by a surface has now provided an alternative way of making a BEC [22]. If we can apply the surface-induced evaporative cooling on an atom chip, this new cooling technique can be potentially used to make a continuous coherent atom source for which a spatial-dependent cooling is required, because the efficiency of the evaporative cooling near the surface is controlled by the atom-surface separation that can be spatially varied in a waveguide.

The principle of surface-induced evaporative cooling is very similar to that of the conventional rf-forced evaporative cooling. In rf-forced evaporative cooling, resonant rf radiation selects hot atoms and transfers them from a trapped spin state to an untrapped spin state. After hot atoms, i.e., energetic atoms, fall out of the trap, the average energy of the remaining atoms decreases after rethermalization through collisions. Similarly, when atoms are brought close to a surface, high-energy atoms hit the surface first because of their broader spatial and momentum distribution in a magnetic trap. Hot atoms are lost from the trap after colliding with the surface, and only colder atoms remain; therefore, the cloud cools down.

In this chapter, I describe an experiment where we bring a cold normal cloud close to the chip surface and investigate the surface-induced cooling effect. It is straightforward to bring an atom cloud close to the surface in our experiment for two reasons.

First, the distance between the atom cloud and the surface in our experiment is less than in a conventional chamber because our atoms are captured and trapped on a chip. The typical distance of the cloud from the surface of the chip is only about $100\ \mu\text{m}$. Second, we can change the distance of the cloud from the surface by simply changing the currents in the wires and the transverse bias field that create the magnetic trap and the waveguide, as described in chapter 2.

In our experiment, we can bring atoms close to either an aluminum nitride (AlN) (the substrate material) or copper (Cu) (the wire material) surface. However, either rf noise from the environment coupled into or thermal noise in the conductor wires can cause reduction of the lifetime [16, 17, 22, 23]. Moreover, the cold atom cloud experiences a corrugated potential due to current noise or imperfection of the wires, such as rough edges or defects. When the variation of the potential is comparable to the energy of the cloud, the cloud breaks up into pieces and is thus fragmented [18, 19]. Therefore, we use the AlN substrate to investigate surface-induced cooling.

We show that with careful design of the trajectory to bring the atoms close to the surface, this cooling technique can efficiently create a BEC from an atom cloud with an initial temperature of $30\ \mu\text{K}$.

3.1 A simple trajectory for bringing atoms close to a surface

We can use a relatively simple trajectory to bring the atom cloud close to the chip surface. Atoms are first trapped above a single wire, which we call the “main wire”. We ramp up or down simultaneously a transverse bias field (B_t) and current of another wire (I_2), which we call the “secondary wire”, to bring the atoms close to the surface. The secondary wire is parallel to the main wire (see Figure 3.1). The change of the current and the transverse bias field shifts the bottom of the trapping potential and moves atoms. The atom cloud is moved to the surface between the two wires by maintaining a constant distance from the main wire so that the cloud simply rotates

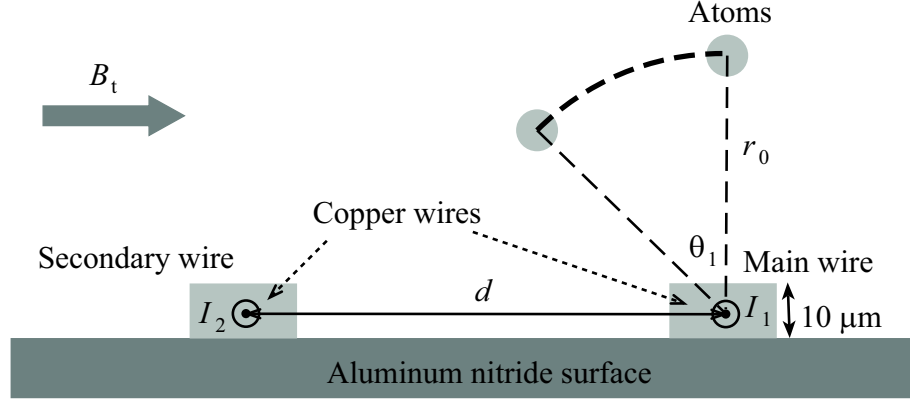


Figure 3.1: The trajectory of the atom cloud. The cloud starts above the center of the main wire and then rotates to the center between the two wires at various distances from the aluminum nitride substrate. θ_1 is the angle of rotation, r_0 is the distance between the atoms and the center of the main wire, d is the distance between the two wires and B_t is the transverse bias field.

from the top to the left side of the main wire, as shown in Figure 3.1.

Because we keep the current of the main wire (I_1) constant, the current in the secondary wire and the transverse bias field can be determined by the angle of the rotation as follows:

$$I_2 = I_1 \frac{\sin \theta_1 (r_0^2 + d^2 - 2dr_0 \sin \theta_1)}{r_0(d - r_0 \sin \theta_1)} \quad (3.1)$$

and

$$B_t = I_1 \frac{\mu_0}{2\pi} \frac{d \cos \theta_1}{r_0(d - r_0 \sin \theta_1)}, \quad (3.2)$$

where θ_1 is the angle of rotation, r_0 is the distance between the atoms and the center of the main wire, d is the distance between the two wires, and μ_0 is the magnetic permeability. When θ_1 is equal to 0° , atoms sit above the main wire with a radial trap frequency of 1.8 kHz. When θ_1 is equal to 90° , the distance of the cloud from the surface is $5 \mu\text{m}$, which is half of the wire thickness, and the radial frequency is 3.6 kHz.

The atom cloud starts from $100 \mu\text{m}$ above the center of the main wire. The initial transverse bias field is 20 G, and the longitudinal bias field is 1 G. The currents in the main wire and the secondary wire are 1 A and 0 A, respectively. The distance between

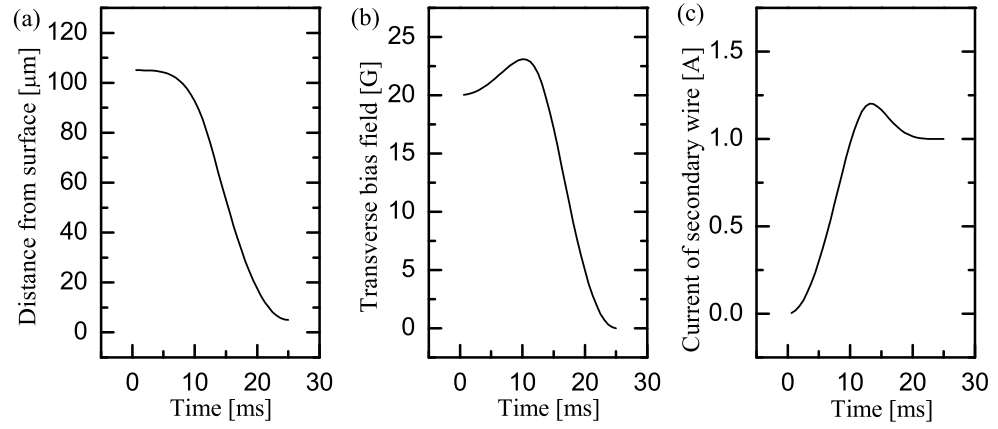


Figure 3.2: (a) The calculated distance, (b) transverse bias field ,and (c) current of the secondary wire as functions of time when the cloud is moved closer to the surface. The cloud is moved closer to the surface in 25 ms with the final position 5 μm above the surface.

the center of the two wires is $200\ \mu\text{m}$. We change the transverse bias field and the current of the secondary wire in 25 ms to move the cloud closer to the surface. We change the cloud position by keeping a constant angular velocity, i.e., $\frac{d\theta_1}{dt} = \text{constant}$. Therefore, Equations 3.1 and 3.2 can be expressed as functions of time (t) by substituting θ_1 with $\text{constant} \times t$. The change of the secondary-wire current and the transverse bias are plotted as functions of time in Figure 3.2. The distances from the surface for each time are calculated by numerically solving for the cloud position, which is at the minimum of the magnetic potential. Figure 3.2 also shows the calculated distance from the surface when we bring the cloud to $5\ \mu\text{m}$ above the surface of the substrate.

We observe the temperature of the cloud by imaging its size (σ) after free expansion. The temperature is proportional to the square of the size, $T \propto \sigma^2$. The change of cloud size as a function of distance from the surface is shown in Figure 3.3. The temperature starts to decrease when the distance from the surface is less than approximately $10\ \mu\text{m}$. Loss of atoms is also observed at the same distance. In contrast to the expected cooling at even shorter distances, the experimental result shows that the cloud is heated rapidly when it comes to around $6\ \mu\text{m}$ from the surface.

To examine the source of heating, we perform a second experiment in which we bring the atom clouds to $5\ \mu\text{m}$ from the surface, hold them at the same distance for various times ranging from 5–7 ms, and then bring them back to their original positions. The change in the number of atoms and cloud sizes for various hold times is shown in Figure 3.4. The cloud heats up for a short hold time, and the cloud temperature reaches a saturation point for a longer hold time. This heating is due to a nonadiabatic change in the radial trap frequency arising from potential minima crossing, and the details are discussed in the following section.

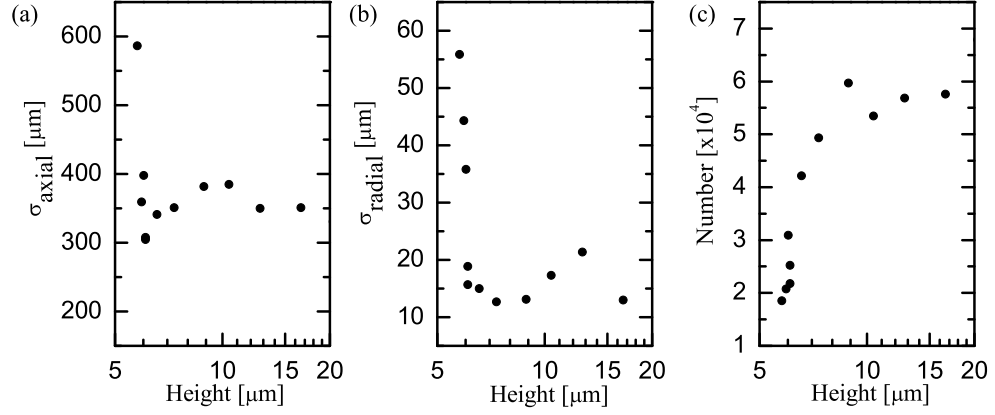


Figure 3.3: (a–b) Cloud sizes and (c) atom numbers at various heights above the surface. Atoms are brought to various heights above the surface in 25 ms and then immediately moved back in another 25 ms to their original position. The cloud size remains constant at longer distances, starts to decrease as the distance decreases, and rapidly increases as the distance approaches 6 μm . The atom number also shows a rapid decline when the distance is within 10 μm from the surface.

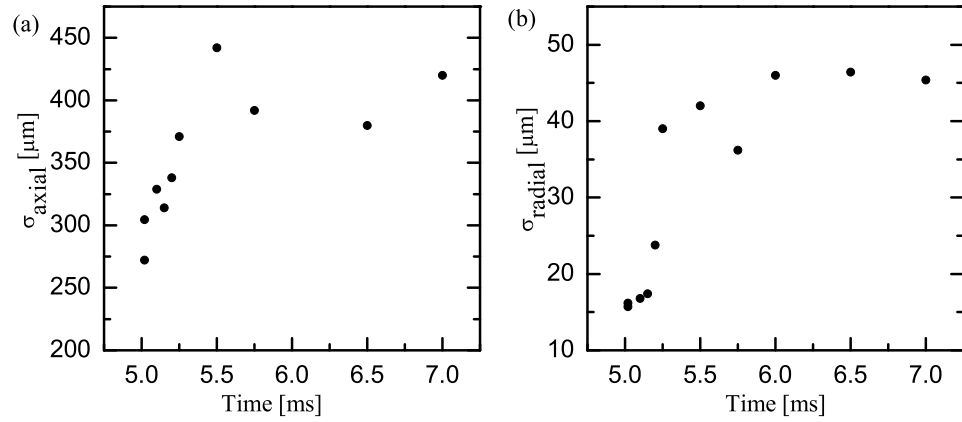


Figure 3.4: Cloud sizes for various hold times close to the surface. Atoms are held 5 μm away from the surface for various time periods. The cloud sizes in both the radial and axial directions increase for a short period of time and then reach a constant value.

3.2 Heating due to potential minima crossing

In order to observe surface-induced cooling, we need to understand the origin of the heating and reduce the heating rate before we can quantitatively study and improve the cooling efficiency. We carefully investigate the current-ramping process by evaluating the change of the radial trap frequency when we bring atoms close to the surface. When we ramp up the current in the secondary wire and ramp down the transverse bias field, atoms follow the magnetic potential minimum, which is initially above the center of the main wire and moves toward the surface. However, there is also a second potential minimum arising from the secondary wire. This second minimum moves toward the center above the two wires, as shown in Figure 3.5, where the (horizontal, vertical) positions of the centers of the main wire and the secondary wire are at $(0 \mu\text{m}, 5 \mu\text{m})$ and $(-200 \mu\text{m}, 5 \mu\text{m})$, respectively. When the atoms move closer to the surface, the two minima get closer. Consequently, the potential barrier between the two minima decreases and the trap becomes loose. The two minima separate again after they cross each other and the trap frequency then increases again. Figure 3.6 shows how the calculated radial trap frequency changes as a function of time as we bring the cloud to $5 \mu\text{m}$ above the surface in 25 ms. The radial trap frequency decreases from around 2 kHz down to 1 kHz and then increases to 4 kHz. The rapid change of the radial trap frequency in such a short time, i.e., 10 ms, causes the cloud to heat up. Because the potential barrier has a minimum value of 1.4 MHz in the crossing region, and this minimum is comparable to the initial energy of the cloud (~ 1.35 MHz), once the atoms heat up, the hot atoms have enough energy to move across this potential barrier. As a result, an atom loss is observed. Besides, ongoing evaporative cooling occurs because the energetic atoms fall out of the trap and only the cold atoms remain. This ongoing evaporative cooling counteracts the heating, and thus the temperature reaches a saturation point.

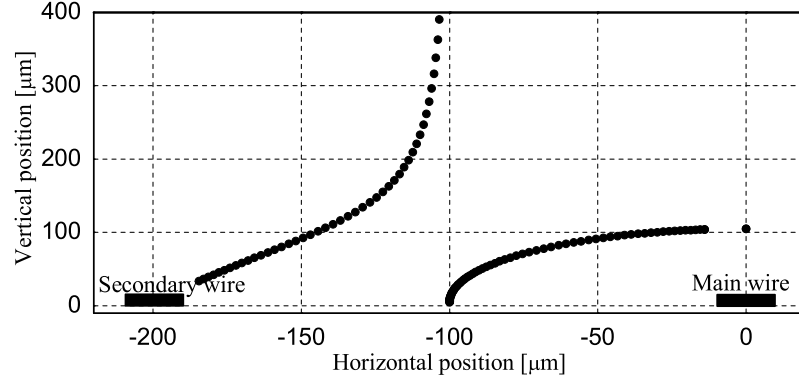


Figure 3.5: Trajectories of the two minima of the magnetic field. The initial position of the cloud at one of the minima is at horizontal position $0 \mu\text{m}$ and vertical position $105 \mu\text{m}$. The two minima come closer and then separate as the cloud moves closer to the surface. The (horizontal, vertical) positions of the centers of the main wire and the secondary wire are at $(0 \mu\text{m}, 5 \mu\text{m})$ and $(-200 \mu\text{m}, 5 \mu\text{m})$, respectively.

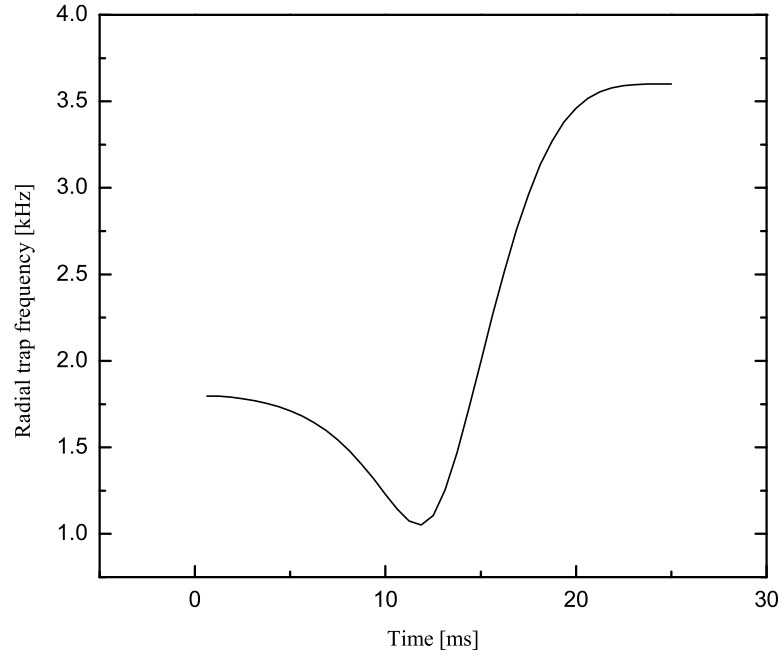


Figure 3.6: The radial trap frequency as a function of time as the cloud is moved closer to the surface. The final position of the cloud is $5 \mu\text{m}$ from the surface, and the total time to bring the cloud to the surface is 25 ms.

3.3 Observation of surface-induced cooling

To reduce the heating of the cloud, we modify the trajectory of the atom cloud from an arc, which has constant distance from the main wire, to a parabola. The parabolic trajectory increases the distance between the two minima when they cross each other and improves the adiabaticity of the ramping process. The relation between horizontal and vertical position of the cloud, X and Y , can be written as follows:

$$Y = A \times \sqrt{X - B}, \quad (3.3)$$

where A and B are two constants that can be determined by the initial and final position. First, by substituting the initial and final (horizontal, vertical) positions, $(0, h)$ and $(-d/2, 0)$, into Equation 3.3, we can solve for A and B , where h is the initial distance between the atoms and the center of the main wire. Second, we substitute $r \sin \theta_1$ and $r \cos \theta_1$ for X and Y in Equation 3.3 and solve for r , the distance between the atoms and the main wire. The distance from the main wire, r , is no longer a constant but rather is a function of the rotation angle, θ_1 , as below:

$$r = \frac{1}{d} h^2 \sec^2 \theta_1 (\sin \theta_1 - \sqrt{\frac{d^2}{h^2} \cos^2 \theta_1 + \sin^2 \theta_1}). \quad (3.4)$$

If we replace the constant r_0 in Equations 3.1 and 3.2 with r in Equation 3.4, we obtain the new trajectory shown in Figure 3.7. In this case, we also change the separation of the two wires from $200 \mu\text{m}$ to $300 \mu\text{m}$ by choosing another wire that is $300 \mu\text{m}$ away from the main wire as a secondary wire. This separation provides a better match of the initial and final radial trap frequencies without changing the current in the main wire. For a 3 G longitudinal bias field, 1 A main-wire current, and a 20 G transverse bias field, the cloud with an initial temperature of $30 \mu\text{K}$ starts initially $100 \mu\text{m}$ above the center of the main wire. As a result, the radial trap frequency remains approximately 1 kHz with a variation of less than 0.3 kHz, as shown in Figure 3.8. After employing the modified trajectory, we observed the size and the atom number of the cloud again by

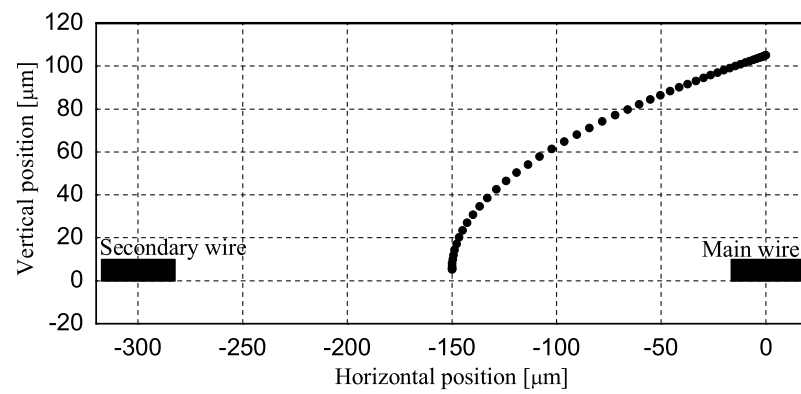


Figure 3.7: A Parabolic trajectory of the cloud for surface-induced cooling. The cloud starts $50\text{ }\mu\text{m}$ above the center of the main wire and then moves in between the two wires.

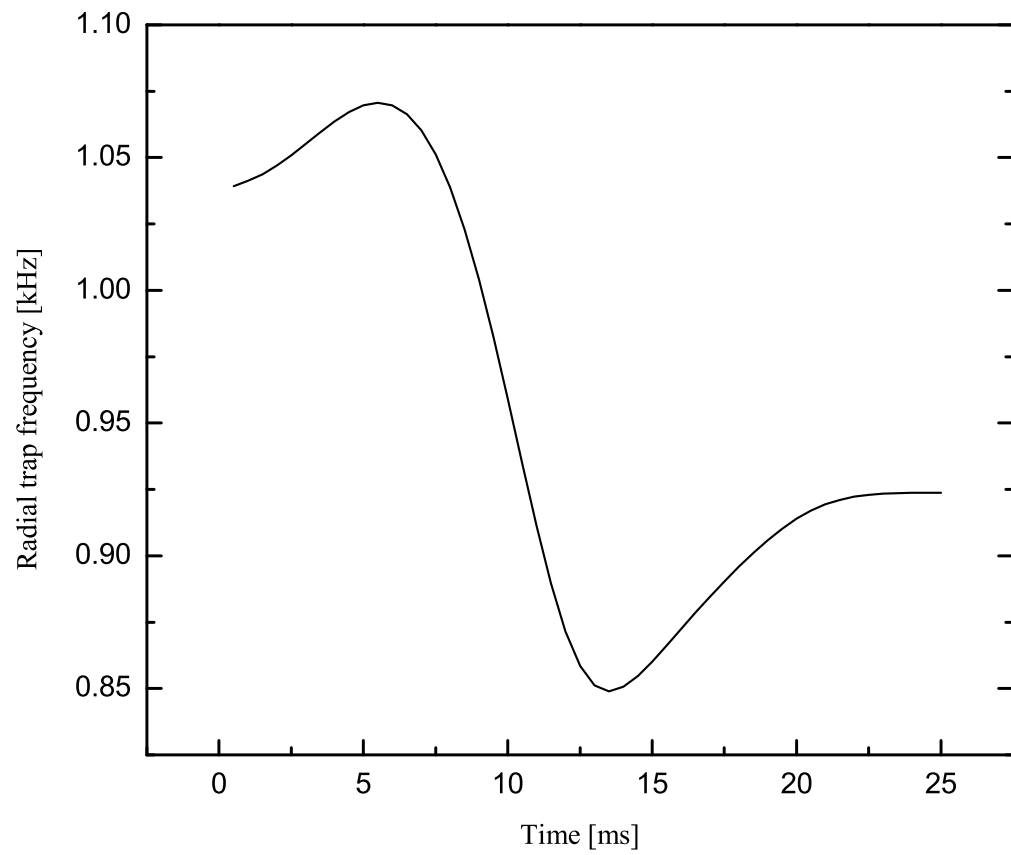


Figure 3.8: The radial trap frequency of a parabolic trajectory as a function of time for surface-induced cooling.

varying the distance from the surface, as shown in Figure 3.9 or the hold time close to the surface, as shown in Figure 3.10. In contrast to the heating observed previously in Figures 3.3 and 3.4, cooling is observed as results of both decreasing the distance from the surface (see Figure 3.9) and increasing the hold time close to the surface (see Figure 3.10).

The atom cloud in a harmonic trap has a spatial Gaussian distribution and more than 84% of the atoms are distributed within the rms size of the cloud. The cooling effect is observable when a sufficient amount of hot atoms, e.g. 10%, are removed from the trap. Thus, we expect the cooling effect occurs when the distance is comparable the rms size of the cloud. For a cloud with an initial temperature of 30 μK in a 1 kHz trap, the in-trap size is around 7.5 μm . In Figure 3.9, the cooling effect appears when the distance from the surface is less than 10 μm , which is close to the expected value. In Figure 3.10, the cooling is also observed as we increase the hold time with the cloud at 7 μm from the surface. A BEC is achieved with a hold time longer than 140 ms.

A BEC is also made by using a different adiabatic trajectory. We change the trajectory of the cloud by moving the cloud to the other side of the main wire, which is farther away from the secondary wire, as shown in Figure 3.11. To increase the adiabaticity of the current ramping process, we increase the radial trap frequency by a factor of two, i.e., from 1 kHz to 2 kHz, by decreasing the current in the main wire from 1 A to 0.5 A. Because the radial trap frequency is proportional to the magnetic gradient ($B' \propto \frac{B_L^2}{r}$) created by the current in the main wire, when we decrease the current to half, the radial frequency is doubled. The current of the secondary wire and the transverse bias field as a function of time can still be determined by Equations 3.1 and 3.2. For a 3 G longitudinal bias field and a 20 G transverse bias field, the cloud is initially located 50 μm above the center of the main wire. The trajectory also maintains the same distance from the main wire, which is located at the origin, as shown in Figure 3.11. Consequently, the radial trap frequency remains approximately 2 kHz with a

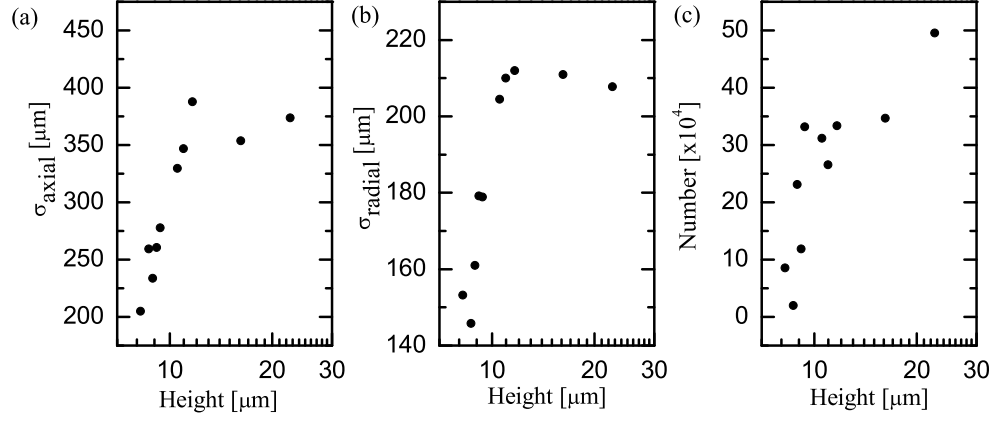


Figure 3.9: Cooling of the atoms as a function of height above the surface. Atoms are brought to various heights from the surface in 25 ms. The cloud size remains constant until the distance is less than 10 μm . Atom loss is observed in the same distance range because of the surface-induced cooling effect.

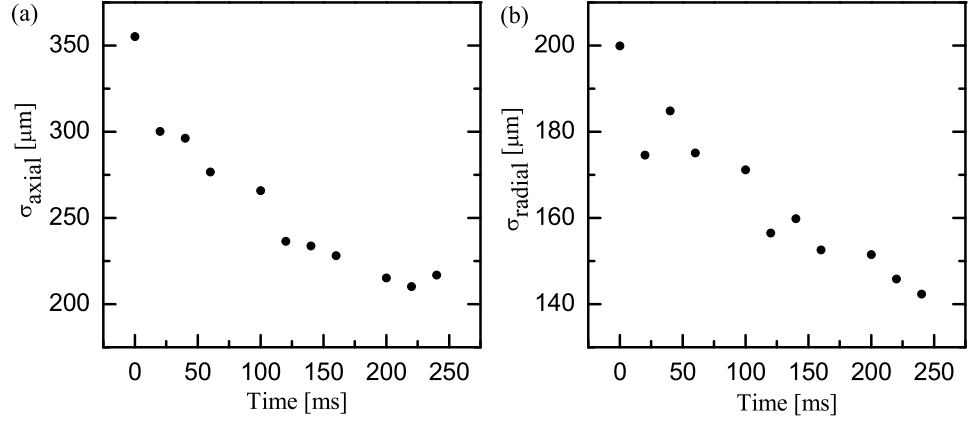


Figure 3.10: Cooling of the atoms for various hold times close to the surface. Atoms are held 7 μm away from the surface for various times. The cloud sizes in both the radial and axial directions decrease as the hold time increases. The free expansion time before taking the images is 10 ms.

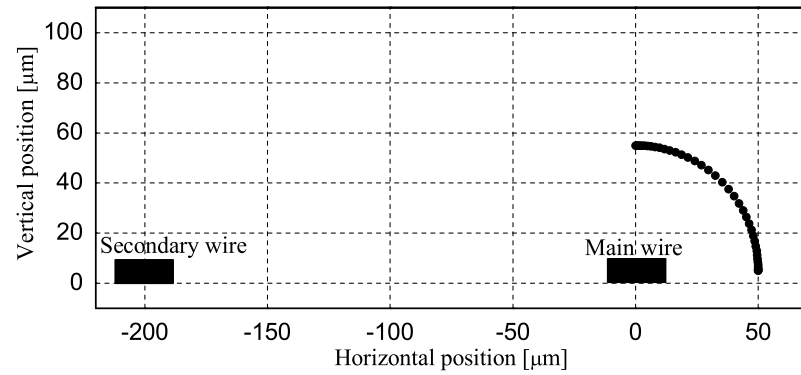


Figure 3.11: An alternative trajectory of the cloud for surface-induced cooling. The cloud starts 50 μm above the center of the main wire and then moves to outside the wires.

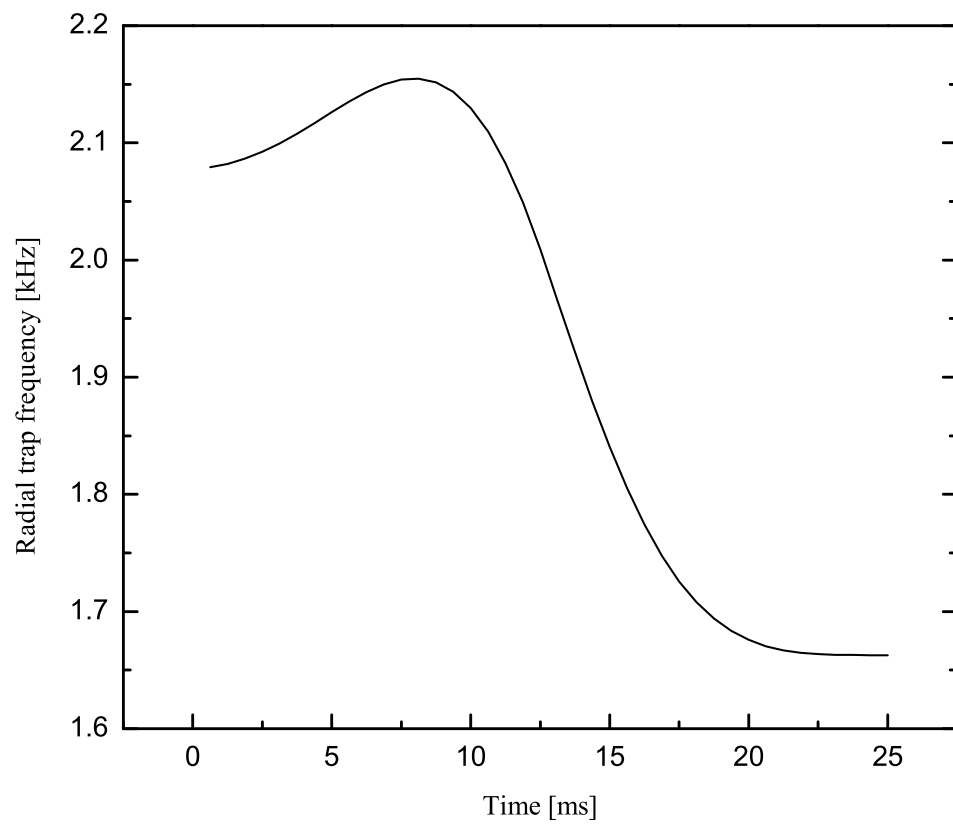


Figure 3.12: The radial trap frequency of the alternative trajectory as a function of time for surface-induced cooling.

variation of less than 0.6 kHz. The variations of the radial trap frequencies by using the adiabatic trajectories introduced in this section are much smaller compared to the relatively simple trajectory introduced in the section 3.1.

We achieve a BEC of around 1.5×10^5 atoms with a cloud at an initial temperature of $30 \mu\text{K}$. The atom number of the condensate formed by using surface-induced cooling is larger than the number of the condensate atoms formed by using only the conventional rf approach, typically $\sim 7\text{--}8 \times 10^4$. While our evaporative cooling still involves rf radiation in the initial cooling stages, the experimental result demonstrated by D. Harber et.al. [22] shows that a BEC can be achieved by using only surface-induced evaporative cooling after the atoms are collected in a MOT and loaded into a magnetic trap. Their result has proved that the surface-induced cooling can be made very efficient and used to replace the conventional rf-forced evaporative cooling technique.

In addition, the surface-induced evaporative cooling has two advantages. First, in contrast of the rf-induced cooling, in which rf radiation only selects the hot atoms within a certain energy bandwidth and leaves other hot atoms trapped (so-called the “Oort cloud”) [24], surface-induced cooling drives out all the hot atoms above a certain energy. Therefore, heating [25, 26, 27, 28, 29] due to the presence of an Oort cloud is reduced. Second, the surface-induced cooling can be well suited for making a cw coherent atom source where a spatial-dependent evaporative cooling is required. It is difficult to apply spatial-varying rf radiation, but it is easy to control the efficiency of the evaporative cooling near the surface by spatially varying the atom-surface separation in a waveguide.

Chapter 4

Diffraction of atoms with a standing-light wave

The diffraction of a matter wave by an optical light wave is analogous to the diffraction of light by a periodic structure, such as a diffraction grating, in optics. It was predicted in 1933 by Kaptiza and Dirac that electrons can be diffracted by a standing light wave [30]. It was pointed out later by Altshuler et al. that the diffraction by standing light waves could also occur for neutral atoms or any other particles that are capable of scattering photons [31]. The diffraction of atoms was first demonstrated in free space with a near-resonant standing wave by Pritchard and his colleagues at MIT [32, 33]. Their observation proved the photon counterpart of matter waves. The same technique were also demonstrated with a coherent wave packet in 1999 [34] when Ovchinnikov et al. successfully diffracted a BEC released from a magnetic trap. Later in 2004, Prentiss and her colleagues at Harvard used an optical standing wave pulse to split a normal cloud longitudinally in a “stadium”-shaped magnetic guide [13]. Therefore, the diffraction of atoms by a standing wave light field has brought up the possibility of a coherent beamsplitter, which is essential for an atom interferometer and a standing light wave is served to split, reflect, and recombine the atoms in our experiment to form an atom interferometer.

In this chapter, we show an experiment where off-resonance laser light is used to create a standing light wave to diffract Rb atoms. After interacting with the standing light wave, atoms are diffracted into different momentum states and propagate with

different velocities. We discuss the diffraction efficiency of atoms theoretically in two different regimes, the Raman-Nath regime and Bragg regime, and quantitatively compare it to our experimental results. The diffraction in different regimes allows us to split or reflect atoms with a high efficiency to make an interferometer. For example, the splitting of the atoms involves a Raman-Nath pulse and the reflection of atoms employs a Bragg process. We first demonstrate the diffraction effect of atoms in both free space and a waveguide using a Raman-Nath pulse. With further control of the light pulse, e.g. by changing intensity and pulse width, we can use this diffraction effect to create a coherent beamsplitter for a BEC. We can also use a Bragg pulse to reflect the atoms in the waveguide so the split cloud could recombine at their origin to form an interferometer.

4.1 Theoretical consideration

To evaluate the diffraction efficiency and understand the diffraction mechanism, we treat the atom motion quantum mechanically. In the electric dipole approximation, the Hamiltonian for an atom in an electromagnetic field with the wave vector in the x-direction takes the form

$$H = H_0 + H_{AF} = H_0 - \vec{\mu} \cdot \vec{E}(x, t), \quad (4.1)$$

where $H_0 = \frac{P_x^2}{2m_{Rb}} + U_0$ is the unperturbed Hamiltonian describing an atom in the absence of light, H_{AF} is the atom-photon interaction term, $\vec{\mu}$ is the dipole moment of the atoms, and $\vec{E}(x, t) = \vec{\epsilon} 2E_0 \cos(kx) \cos(\omega t)$ is the standing-wave field with peak amplitude $2E_0$, wave number k , frequency ω , and polarization $\vec{\epsilon}$. A standard method used to describe and understand this coupled atom-photon system is the so called “dressed state approach” [35, 36], where the laser field is treated as a quantized field rather than a classical external field. We transfer the basis into a rotating frame, the rotating frequency of which is equal to the frequency difference (ω_0) between the unperturbed

excited and ground states, and introduce a Rabi frequency

$$\Omega_0 = \frac{\vec{\mu} \cdot \vec{\epsilon} E_0}{\hbar}, \quad (4.2)$$

where the Rabi frequency indicates the strength of the coupling between the ground and excited states. The coupling is stronger when the electric field is stronger and thus the Rabi frequency is larger. In most situations, the laser fields are near resonant with the atomic transitions, and it is an excellent approximation to ignore the rapidly oscillating, or counterrotating terms. This rotating-wave approximation yields for the interaction Hamiltonian a simpler form

$$H'_{AF} = \Omega_0 \cos kx = \frac{\Omega_0}{2}(e^{ikx} + e^{-ikx}). \quad (4.3)$$

After scattered by the light field, atoms can be either in the electronic excited or ground state, so the wave function $\psi(x, t)$ can be expressed as a superposition of an excited and ground electronic state, $|e\rangle$ and $|g\rangle$:

$$|\psi(x, t)\rangle = e(x, t) |e\rangle + g(x, t) |g\rangle, \quad (4.4)$$

where $e(x, t)$ and $g(x, t)$ are the coefficients as functions of position (x) and time (t). Substituting the wave function into Schrödinger's equation, $i\hbar \frac{\partial \psi}{\partial t} = H\psi$, yields the coupled equations:

$$i\hbar \frac{\partial e(x, t)}{\partial t} = -\frac{\hbar^2}{2m_{Rb}} \frac{\partial^2}{\partial x^2} e(x, t) + \hbar\Omega_0 \cos(kx)g(x, t) - \hbar\Delta\omega e(x, t), \quad (4.5)$$

and

$$i\hbar \frac{\partial g(x, t)}{\partial t} = -\frac{\hbar^2}{2m_{Rb}} \frac{\partial^2}{\partial x^2} g(x, t) + \hbar\Omega_0 \cos(kx)e(x, t), \quad (4.6)$$

where $\Delta\omega$ is the detuning of the light from the atomic resonances. If we know the intensity of the laser beam, i.e., the Rabi frequency, and the interaction time, we can solve the coupled equations for the probabilities of the different states. We can further determine diffraction efficiencies for different momentum states by taking the Fourier transform of the wave functions.

The diffraction effects are separated into two different regimes: Raman-Nath and Bragg regime. In the Raman-Nath regime, the uncertainty of the energy associated with the interaction time is larger than the energy splitting between the different momentum states and the energy of the atom-photon system is not conserved in the scattering process. In contrast, in the Bragg regime, the uncertainty of the energy associated with the interaction time is smaller than the energy splitting between the different momentum states, and the energy of the atom-photon system is conserved.

4.1.1 Raman-Nath regime

If the atom cloud is initially at rest, a short pulse enables the diffraction process by giving atoms both a momentum kick and additional kinetic energy. In the Raman-Nath regime, the atom-photon interaction time is so short that the Doppler shift of the resonance frequency due to atom recoil is negligible. In general, the resonance frequency will have a Doppler shift $\Delta\omega_{Doppler} \propto kp/m_{Rb}$ if the momentum of atoms (p) is nonzero, so the resonance frequencies of atoms in different momentum states are slightly different. In the Raman-Nath regime, however, if the atom-photon interaction time is so short that the frequency width due to the uncertainty arising from finite interaction time is much larger than the Doppler shift, this Doppler frequency shift can be ignored. In practice, this means that pulse lengths less than about $133 \mu s$ are needed for Rb atoms. Since the recoil energy of an atom, $E_{rec} = \hbar\omega_{rec} = \frac{\hbar^2 k^2}{2m_{Rb}}$, is half the Doppler shift energy, $E_{doppler} = \hbar\omega_{Doppler} = \frac{\hbar kp}{m_{Rb}}$, the Raman-Nath regime also implies that the interaction time is short compared to the inverse recoil frequency. Once we ignore the Doppler shift, the frequency of the standing-light wave can be considered as on resonance for all the momentum states. The kinetic energy terms, $-\frac{\hbar^2}{2m_{Rb}} \frac{\partial e(x,t)}{\partial x^2}$ and $-\frac{\hbar^2}{2m_{Rb}} \frac{\partial g(x,t)}{\partial x^2}$, are also negligible based on the same argument. Equations 4.5 and 4.6

can then be rewritten as

$$i\hbar \frac{\partial e(x, t)}{\partial t} = \hbar\Omega_0 \cos(kx)g(x, t) - \hbar\Delta\omega e(x, t), \quad (4.7)$$

and

$$i\hbar \frac{\partial g(x, t)}{\partial t} = \hbar\Omega_0 \cos(kx)e(x, t), \quad (4.8)$$

4.1.1.1 On-resonance standing light wave

If we first consider the simplest case where the standing light wave is on resonance with an atomic transition, i.e, $\Delta\omega = 0$, the last term in Equation 4.7 vanishes. Because momentum is conserved in the process and photon momentum is quantized, atom momentum can only be increased by multiples of the photon momentum, $\hbar k$. Therefore, the partial wave function $e(x, t)$ and $g(x, t)$ can be expanded as

$$e(x, t) = \sum_m e_m(t) e^{i(\frac{p_i}{\hbar} + mk)x} \quad (4.9)$$

and

$$g(x, t) = \sum_m g_m(t) e^{i(\frac{p_i}{\hbar} + mk)x}, \quad (4.10)$$

where p_i is the initial momentum of the atoms and $e_m(t)$ and $g_m(t)$ give the probabilities of diffraction into the m th diffraction orders. For a given initial condition in which atoms start in the ground state with center of mass motion at rest, $e_m(t) = 0$, $g_m(t) = \delta_{m0}$, and $p_i = 0$, Equations 4.9 and 4.10 simplify to

$$i\hbar \frac{dc_m(t)}{dt} = \frac{\hbar\Omega_0}{2} [c_{m-1}(t) + c_{m+1}(t)], \quad (4.11)$$

where

$$c_m(t) = \begin{cases} e_m(t) & m = \text{odd} \\ g_m(t) & m = \text{even} \end{cases}. \quad (4.12)$$

The solutions of this equation are the Bessel functions of the first kind [37] so we have

$$c_m(t) = i^m J_m(\Omega_0 t), \quad (4.13)$$

and the probability of atoms in one particular momentum state is

$$p_m(t) = J_m^2(\Omega_0 t), \quad (4.14)$$

As shown in Equation 4.14, the probability of atoms in the m th diffraction order can be evaluated as a Bessel function of order m and depends on standing wave intensity and atom-photon interaction time. Because Bessel functions have the property that $J_{-m}(x) = (-1)^m J_m(x)$, we thus have the probability of the m th order equal to that of the $-m$ th order and the atom number distributions in the positive and the negative order are symmetric. Because the atoms are initially in the ground state, they are diffracted into an electronic excited state after scattering an odd number of photons, i.e., when m is an odd number, and a ground state after scattering an even number of photons, i.e., when m is an even number.

As a result of diffraction, atoms are transferred to different momentum states. However, this diffraction mechanism with resonant light is not a coherent process. If an atom absorbs an odd number of photons, it jumps to an excited state. An atom that is in an excited state will spontaneously emit a photon in an arbitrary direction as it falls back to the ground state. Thus the atom gains a photon momentum ($\hbar k$) in a random direction. Consequently, the atom momentum distribution is broad due to spontaneous emissions.

4.1.1.2 Off-resonance standing light wave

The spontaneous emissions increase the dispersion of the atom momentum and therefore, destroy the coherence of the wave packets. To reduce this decoherence effect, we can use an optical standing wave with a frequency detuned from the atomic resonance to avoid significantly populating the excited state of the atoms. We can expand the theoretical model to include an off-resonance photon-scattering process. For a large

detuning

$$|\Delta\omega| = |\omega - \omega_0| \gg \Omega_0, \omega_{rec}, \quad (4.15)$$

where $\omega_{rec} = \frac{\hbar k^2}{2m_{Rb}}$ (recoil frequency). Atoms that are initially in their ground state, will have a negligible probability of being in the excited state after diffraction. For $\Delta\omega < 0$, Equation 4.6 can also be approximated [35, 36, 38, 34] as

$$\begin{aligned} i\hbar \frac{\partial g(x, t)}{\partial t} &= -\frac{\hbar^2}{2m_{Rb}} \frac{\partial^2}{\partial x^2} g(x, t) + \frac{\hbar^2 \Omega_0^2}{|\Delta\omega|} \cos(kx)^2 g(x, t) \\ &= -\frac{\hbar^2}{2m_{Rb}} \frac{\partial^2}{\partial x^2} g(x, t) + 2\hbar\Omega_{eff} \cos(kx)^2 g(x, t), \end{aligned} \quad (4.16)$$

where an effective Rabi frequency is defined as $\Omega_{eff} = \frac{\Omega_0^2}{2|\Delta\omega|}$. By substituting the wave function in Equation 4.10 into Equation 4.16 and ignoring the kinetic energy, the first term on the right side of Equation 4.16, due to the short interaction time, we get the differential equation

$$\begin{aligned} i\hbar \frac{\partial g_m(t)}{\partial t} &= \left[-\frac{\hbar^2}{2m_{Rb}} (p_i + m\hbar k)^2 + \hbar\Omega_{eff} \right] g_m(t) \\ &\quad + \frac{\hbar\Omega_{eff}}{2} [g_{m+2}(t) + g_{m-2}(t)] \end{aligned} \quad (4.17)$$

$$\simeq \frac{\hbar\Omega_{eff}}{2} (g_{m+2}(t) + g_{m-2}(t)), \quad (4.18)$$

where the term of $\hbar\Omega_{eff}g_m(t)$ is also ignored because it represents a constant phase evolution, which is the same for all different momentum states. By comparing this equation to Equation 4.4, we find the solution in the off-resonance case

$$g_m(t) = i^{\frac{m}{2}} J_{\frac{m}{2}}(\Omega_{eff}t), \quad (4.19)$$

with the probability of atoms in one particular momentum being

$$p_m(t) = J_{\frac{m}{2}}^2(\Omega_{eff}t). \quad (4.20)$$

The diffraction efficiency in an off-resonance case also has the same Bessel function behavior except that the diffraction is now a coherent process because the spontaneous emission is significantly suppressed due to the large detuning.

4.1.2 Bragg regime

For an atom cloud that propagates with a nonzero momentum after being diffracted by a Raman-Nath pulse, the cloud can also be reflected by another optical standing-wave pulse. The reflecting pulse reverses the propagation direction of the atoms without changing their speed. The reflection of the atoms can be achieved by a diffracting pulse in the Bragg regime with a high efficiency, i.e., nearly 100%. In contrast to the Raman-Nath regime, the Bragg regime refers to the case where the atom-photon interacting time is comparable to or longer than the inverse of the recoil frequency, ω_{rec} . The frequency shift due to the different momenta can no longer be ignored and energy conservation must be obeyed in the process. If we examine the kinetic energy before and after the photon scattering process, the change of the kinetic energy (ΔE_k) is

$$\Delta E_k(m) = \frac{1}{2m_{Rb}}[(p_i + m\hbar k)^2 - p_i^2]. \quad (4.21)$$

The only condition that satisfies $\Delta E_k(m) = 0$ is $p_i = -\frac{m}{2}\hbar k$, where m is a multiple of 2. Atoms maintain their speed but change their direction in the diffraction process, i.e., momentum of the atoms changes from $-\frac{m}{2}\hbar k$ to $\frac{m}{2}\hbar k$, with integer $\frac{m}{2}$. In a m th-order Bragg scattering, the conservation of momentum implies that atoms undergo a $2m$ -photon scattering process and go through $2m-1$ intermediate states (see Figure 4.1). For a first-order Bragg diffraction with a detuned standing wave and an initial momentum of $\hbar k$, the coupled equations between $m = \pm 1$ states are

$$i\hbar \frac{\partial g_1(t)}{\partial t} = (\hbar\omega_{rec} + \Omega_{eff})g_1(t) + \frac{\hbar\Omega_{eff}}{2}g_{-1}(t) \quad (4.22)$$

$$i\hbar \frac{\partial g_{-1}(t)}{\partial t} = (\hbar\omega_{rec} + \Omega_{eff})g_{-1}(t) + \frac{\hbar\Omega_{eff}}{2}g_1(t), \quad (4.23)$$

For given initial conditions, $g_1(0) = 1$ and $g_{-1}(0) = 0$, i.e., atoms initially populate the momentum state $|p = \hbar k\rangle$, the probabilities in each momentum state are

$$p_1(t) = g_1(t)^2 = \cos^2\left(\frac{\Omega_{eff}t}{2}\right) \quad (4.24)$$

$$p_{-1}(t) = g_{-1}(t)^2 = \sin^2\left(\frac{\Omega_{eff}t}{2}\right) \quad (4.25)$$

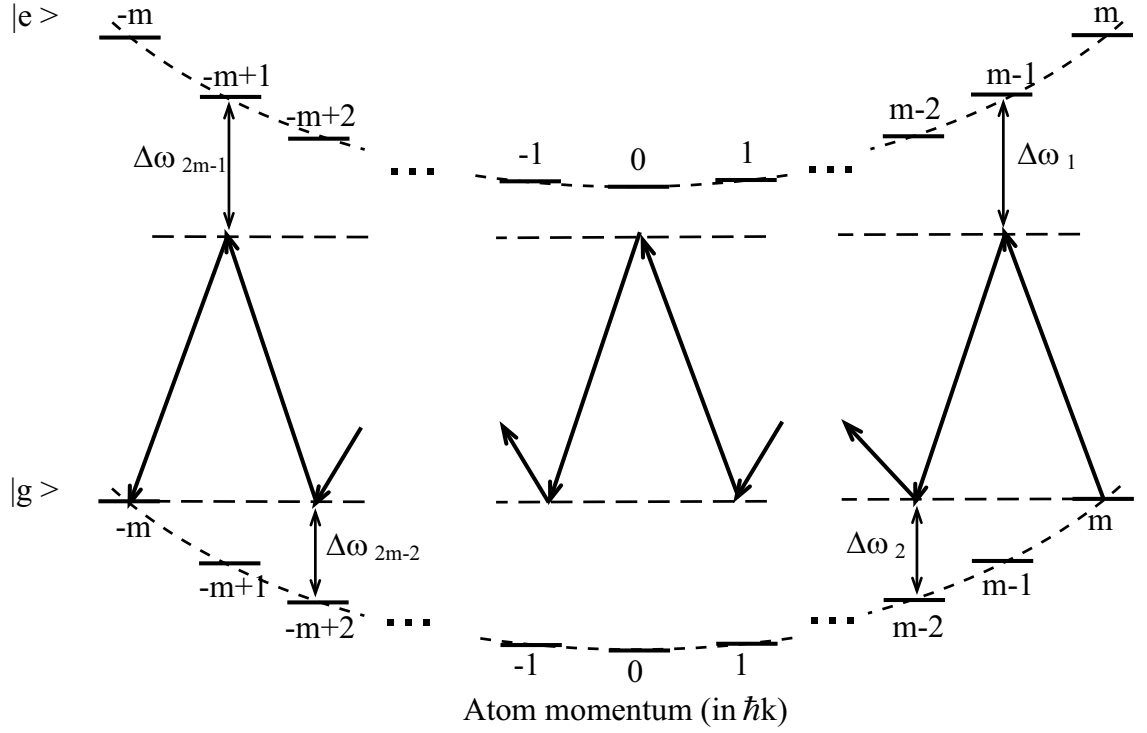


Figure 4.1: Transition diagram for a m th-order Bragg scattering. Atoms start in a state with momentum of $m\hbar k$ and go through $2m-1$ intermediate states in a $2m$ -photon Bragg-scattering process. $\Delta\omega_i$ is the detuning from the i th intermediate state with $i = 1, 2, \dots, 2m - 1$. The vertical axis in this diagram displays the total energy of the atoms, which consists of the internal energy difference between the electronic excited and ground states, and the kinetic energy, E_k . Because $E_k = \frac{\hbar^2 k^2}{2m_{Rb}}$, the different momentum states for a fixed internal energy lie on a parabola in this plot.

The equations can also be expanded to m th-order Bragg diffraction [39, 40] and the probability for $p = \pm m\hbar k$ states are approximated as

$$p_m(t) = g_m(t)^2 = \cos^2\left(\frac{\Omega_m t}{2}\right) \quad (4.26)$$

$$p_{-m}(t) = g_{-m}(t)^2 = \sin^2\left(\frac{\Omega_m t}{2}\right), \quad (4.27)$$

with 2m-photon Rabi frequency

$$\Omega_m = \frac{\Omega_0^{2m}}{2^{m-1}|\Delta\omega_1\Delta\omega_2\ldots\Delta\omega_{2m-1}|}, \quad (4.28)$$

where $\Delta\omega_i$ is the detuning from the i th intermediate level (see Figure 4.1), with $i = 1, 2, \dots, 2m - 1$. If the atoms initially populate the $|p = m\hbar k\rangle$ state, the populations start to transfer to the $|p = -m\hbar k\rangle$ state after the standing-wave light pulse is turned on. After a time $t = \frac{\pi}{2\Omega_{eff}}$, the atoms equally populate both states and the Bragg pulse is called a “ $\frac{\pi}{2}$ ” pulse. After a time $t = \frac{\pi}{\Omega_{eff}}$, all the atoms are transferred to the $|p = -m\hbar k\rangle$ state, i.e., the population in the $|p = m\hbar k\rangle$ state is zero, and the Bragg pulse is called a “ π ” pulse. For a time $\frac{2\pi}{\Omega_{eff}} > t > \frac{\pi}{\Omega_{eff}}$, the population of the atoms in the $|p = -m\hbar k\rangle$ state would start to decrease and the population of the atoms in the $|p = m\hbar k\rangle$ state would thus increase. The frequency with which atoms hop between two states is the effective Rabi frequency, Ω_{eff} , as defined in Equation 4.6.

In our experiment, we use a standing light wave with a detuning ($\Delta\omega$) of -6.8 GHz from the D2 line (the $5^2S_{1/2} \rightarrow 5^2P_{3/2}$ transition) to minimize the spontaneous emission because the spontaneous emission will destroy the coherence of the condensate. The spontaneous emission rate can be estimated as

$$\gamma_{scat} = \frac{\frac{\gamma}{2}s_0}{1 + s_0 + [2(\Delta\omega + \omega_{rec})/\gamma]^2}, \quad (4.29)$$

where γ is the decay rate of the excited state, $s_0 = 2\Omega_0^2/\gamma^2$ is the on-resonance saturation parameter. γ is $2\pi \cdot 6.1$ MHz and the recoil frequency (ω_{rec}) is $2\pi \cdot 3.77$ kHz, for ^{87}Rb D2 line [41]. The recoil energy is velocity dependent and negligible when the detuning

is large. Because the single-photon Rabi frequency in our experiment is $\Omega_0 = 8\text{--}80$ MHz, the spontaneous emission rate is 4.5–450 Hz, as given by Equation 4.29. The spontaneous emission rate is negligible during the time when the standing wave is turned on, typically 1–100 μs .

Our detuning is much larger than both the Rabi frequency and the recoil frequency such that our experimental condition satisfies the criteria described in Equation 4.15, and thus our atomic diffraction probabilities by the standing light wave in the Raman-Nath and Bragg regime are well described by using Equations 4.20 and 4.27.

4.2 Diffraction of atoms in free space

Our first experiment is to observe diffraction of a BEC in a Raman-Nath regime in free space. Our chip is mounted upside down and atoms are actually trapped below the chip surface. When we release the condensate, it falls due to gravity. When the condensate moves close to the center of the optical beams, we turn on an optical standing wave to diffract the atoms. After diffraction, the cloud splits in the direction along the standing wave.

4.2.1 Experimental setup

The diffracting beam has a linear polarization and must be turned on and off quickly to produce a short pulse typically around 1–100 μs . To accomplish this, the diffracting beam is focused into an acousto-optic modulator to control the timing of the pulses. The first order output of the modulator goes through an aperture and reaches the atom chip while the zeroth order is blocked by the aperture (see Figure 4.2). The frequency shift due to the acousto-optic modulator is much smaller than the detuning so this frequency shift can be ignored. We adjust the position of the lens (L4) after the acousto-optic modulator to let the diffracting beam be slowly focused with a waist of 0.6 mm at the retroreflecting mirror (M3). After retroreflection, the diffracting beam

forms a standing wave. Ideally, the imaging beam would be set up perpendicular to the diffracting beam in order to observe the maximum separation of the split clouds. However, our limited optical access forces the diffracting beam and the imaging beam to be less than 90° . Both beams go through the same vacuum window with a angle 2φ between them (see Figure 4.2). In this case, the clouds separation in the images (Δx) and the separation in real space (Δs) are related by

$$\Delta x = \sin 2\varphi \Delta s. \quad (4.30)$$

Because the distance from the center of the chip to the center of window (L) is 7 mm, and the radius of the vacuum window (W) is 2mm, the angle 2φ is given as:

$$2\varphi = 2 \tan^{-1}(W/L) = 2 \times 16^\circ = 32^\circ. \quad (4.31)$$

The imaging beam passes through a pair of lenses, with focal lengths of $f_1 = 100$ and $f_2 = 600$ mm, and goes to the CCD (charge-coupled device) camera.

4.2.2 The alignment of the diffracting beam in free space

The alignment of the diffracting beam is very sensitive because the BEC drops fast in free space and the diffracting beam needs to be centered at the cloud while it is turned on to diffract the atoms. After setting up the standing wave, we have to verify that the diffracting beam is aligned with the condensate and the intensity of the incident and reflected beam are balanced. We change the frequency detuning of the standing wave and block the retroreflected beam. The photon-scattering rate can be described using Equation 4.29, which has a maximum value when the frequency is on resonance, i.e., $\Delta\omega = 0$. When atoms absorb photons, they feel a force ($F = \hbar k\gamma$) in the direction of the incident beam so that the cloud is pushed by the beam. The scattering rate is proportional to the intensity. If the pulse width is very short and the atoms do not move very much when the pulse is turned on, the intensity that atoms experience is nearly

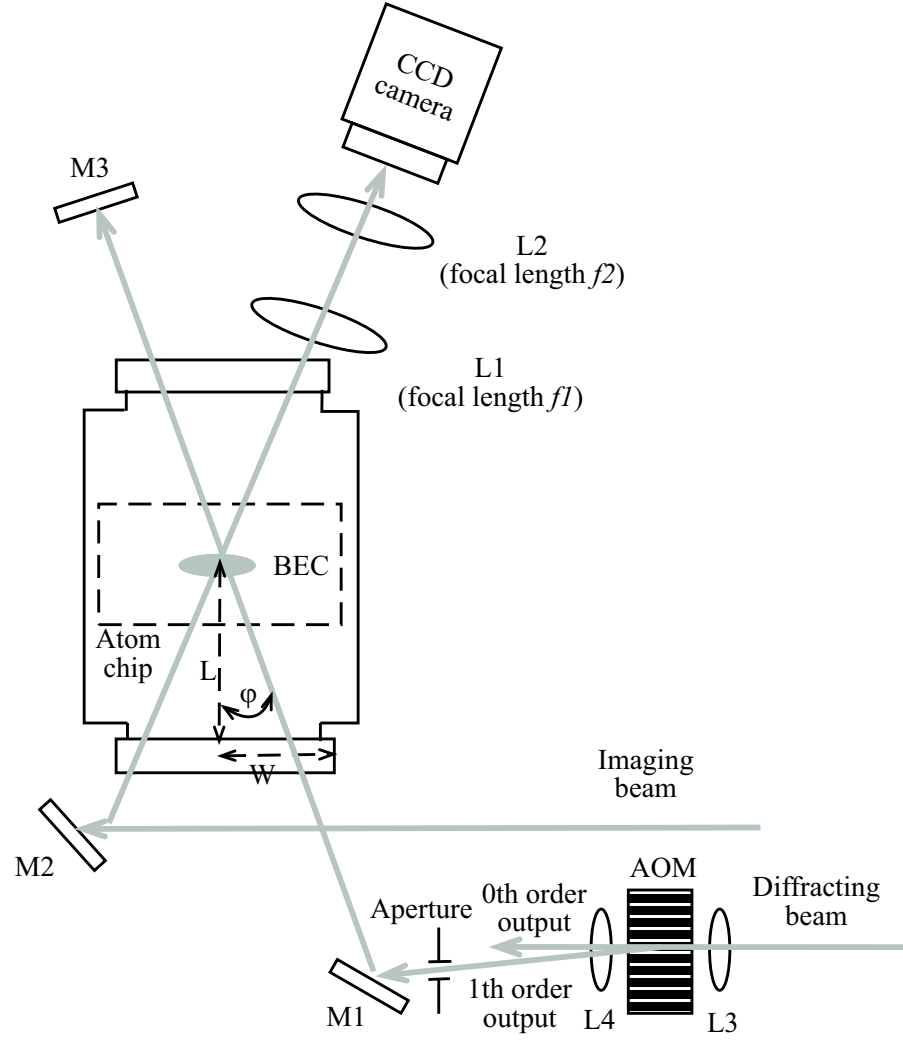


Figure 4.2: Setup of the imaging system and standing wave in free space. The angle φ between the imaging and diffracting beams is 32° . The diffracting beam is focused at an acousto-optic modulator (AOM), which is used to control the timing of the standing wave pulse. The first-order output from the acousto-optic modulator goes to the chip and is retroreflected by a mirror, M3, while the zeroth order is blocked by the aperture.

a constant and the force can be considered constant. The horizontal displacement of the center of the cloud due to the optical force is proportional to square of the atom-photon interaction time, t^2 . After the pulse is turned off, atoms continue to move apart with a constant velocity, $v = \frac{Ft}{m_{Rb}}$, before we take an image. The total time of flight (t_{TOF}) between when the pulse is turned on and when the image is taken is constant, so that the total displacement is $D = \frac{F}{m_{Rb}}(t \cdot t_{TOF} - \frac{1}{2}t^2)$. We observed this horizontal displacement by increasing the interaction time, i.e. the pulse width of the standing light wave. The experimental result is shown in Figure 4.3. We see the displacement of the center of the cloud is a parabolic function of the interaction time, which agrees with the expectation. We also observed heating as an increase of the axial size of the cloud. After absorption, atoms also spontaneously emit photons in random directions. Consequently, the broadening in the momentum distribution causes the increase of the cloud size after expansion, i.e., the atoms heat up.

If we unblock the reflected beam, atoms feel another optical force in the opposite direction and the forces from the incident and reflected beams cancel out. In the experiment, we see that the displacement of the cloud position is reduced as shown in Figure 4.4. However, a nonzero displacement is still visible, particularly for longer interaction times, and implies that the forces are not perfectly balanced. Because the two beams have nearly the same size at the location of the cloud, this result indicates that the power of the reflected beams is weaker than the power of the incident one. The reflected power is attenuated by the vacuum windows. The transmission of each vacuum window is typically around 90–95%. Because the reflected beam passes through the window twice more than the incident beam before it gets to the atoms, the power of the reflected beam is only 80–90% of the power of the incident one. When the two beams are not balanced, the contrast ratio of the standing wave pattern, i.e., the periodical potential, is reduced so that the diffraction efficiency decreases.

We can make the reflected beam size smaller to compensate for the power atten-

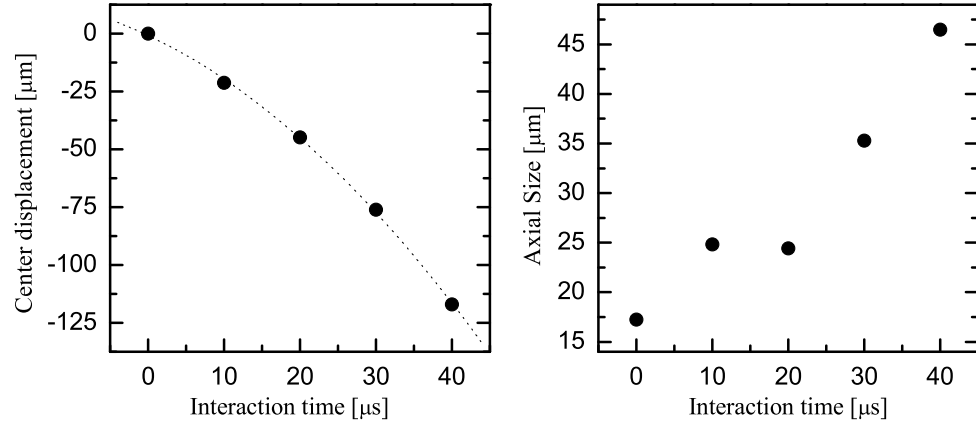


Figure 4.3: Displacements of the cloud centers and changes of the cloud size induced by photon absorption. Left: The position of the cloud center plotted is a parabolic function of the interaction time. The solid circles are the experimental data and the dotted line is a parabolic fitting curve. Right: The heating due to spontaneous emission is observed as an increase of the axial size of the cloud as interaction time increases.

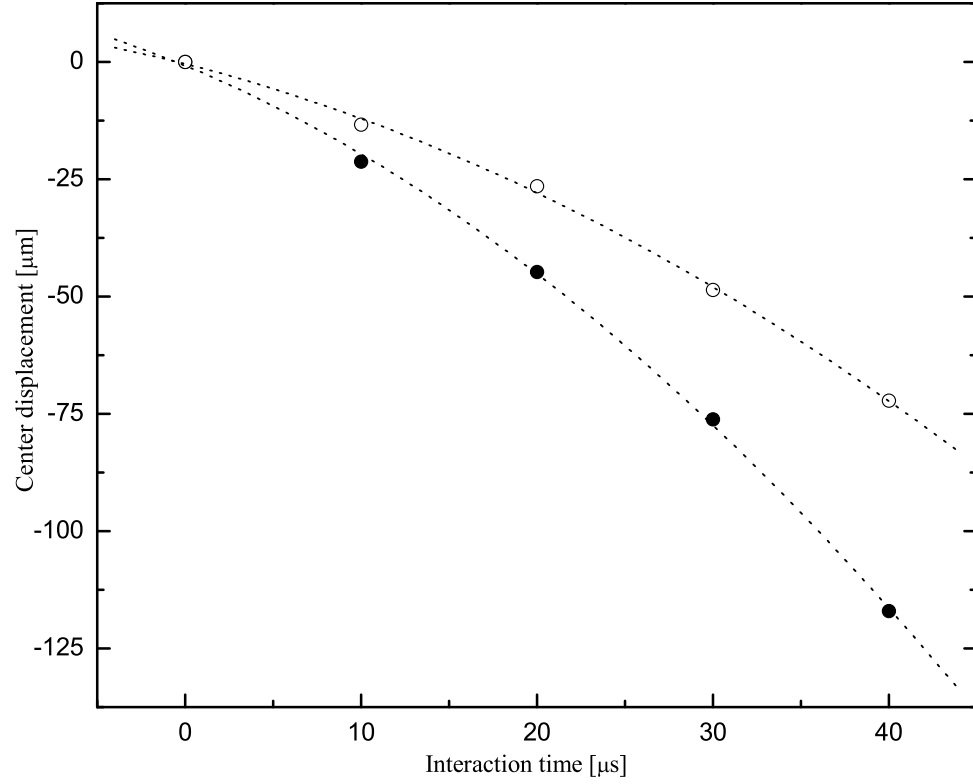


Figure 4.4: Displacements of the cloud centers when the retroreflected beam is on. The solid circles are the experimental data points when only the incident beam is on and open circles are the data points when the incident and retroreflected beams are both on. The displacement is reduced because the optical forces from the two beams partially cancel out.

uation and equalize the intensity of the two beams by moving the lens, L3, in Figure 4.2. We move the lens closer to the acousto-optic modulator to increase the focal length so that the reflected beam is slightly focused after the retroreflecting mirror M3. The results for the two balanced beams are shown in Figure 4.5. The open circles are the data with both beams on after we try to balance their intensity and the solid circles are the data for the single beam. The displacement with two counterpropagating beams is much smaller than the one with single beams and has only a very small change as the interaction time increases. When the intensities of the two beams are nearly equal, the optical forces in the opposite directions cancel so the cloud center doesn't move as we increase the interaction time. We also see that the heating due to spontaneous emission is significantly suppressed when we turn on both beams. If the frequency of the light is not exactly on resonance, the atoms feel a velocity-dependent force ($F' = \hbar k \gamma_{\text{scat}}$) due to the presence of the recoil energy in the scattering rate and this velocity-dependent force provides a one-dimensional (1D) Optical Molasses (OM) cooling [42, 43, 44, 45, 46] for the cloud, which overcomes the heating.

4.2.3 Diffraction patterns

Once the standing light wave is centered at the condensate, we change the effective Rabi frequency or the interaction time of the diffracting beam to vary the diffraction pattern and compare the observed patterns to the theoretical prediction described in Equation 4.20.

We change the frequency of the diffracting beam back to -6.8 GHz detuned from the atomic resonance to avoid spontaneous emissions. The diffracting beam is turned on 7 ms after the BEC is released from the trap and we image 8 ms after turning on the diffracting beam to let the clouds propagate and separate.

The first experiment is done by varying the pulse width of the standing wave. For an input power (P) of 0.61 mW, the output of diffraction is shown in the absorption

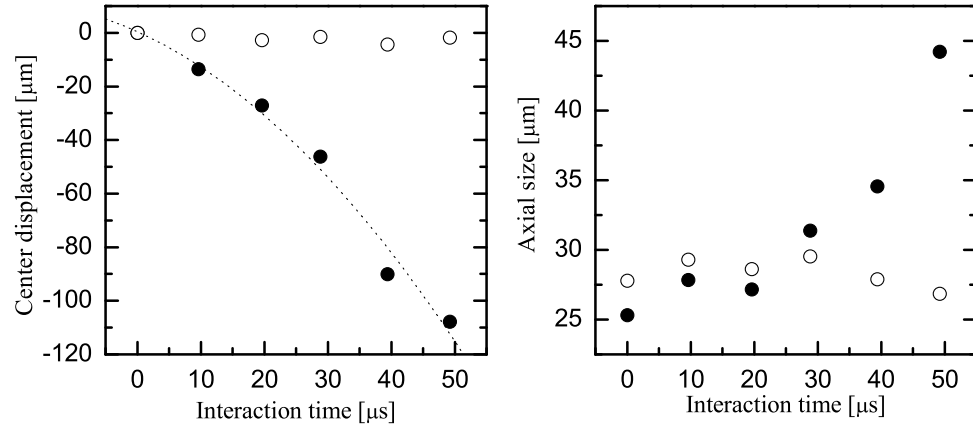


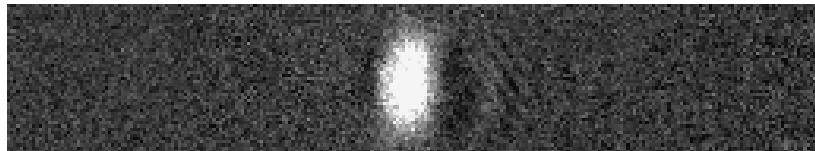
Figure 4.5: Displacements of the cloud centers and axial cloud sizes when the incident and retroreflected beams are balanced. The solid circles are the experimental data when only the incident beam is on and open circles mark the data points when the incident and retroreflected beams are on. Because the intensity of the two beams are equal, the displacement stays nearly zero as we increase the interaction time. The heating due to spontaneous emission is also suppressed because of one-dimensional Optical Molasses (OM) cooling when the beams are both turned on.

images in Figure 4.6 as the interaction time increases. In those absorption images, white represents a high density distribution and black is the background where the density of atoms is zero. We observe a single condensate at $t=0$. As we increase t , the clouds start to populate higher momentum states. The clouds are imaged after 8 ms free propagation such that the atoms in the different momentum states are spatially separated. The separation between the two second diffraction order clouds is measured as $185 \mu\text{m}$. This separation is close to the $199\text{-}\mu\text{m}$ separation calculated from Equation 4.30.

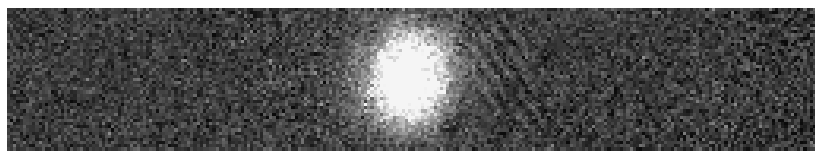
In this case, it is hard to evaluate the atom number and we couldn't quantitatively compare the experimental result with the theoretical model because the separations between the clouds are not large enough. A larger separation can usually be achieved by increasing the propagation time. However, atoms fall rapidly due to the gravitational force and they fall out of the imaging window before the separation becomes large enough to distinguish different momentum components. Instead of evaluating atom number in each component, we numerically generate a diffraction pattern based on Equation 4.20 by assuming the cloud size for all the diffraction orders are the same as the initial size of the cloud and the numbers in each diffraction order are Bessel functions as described in Equation 4.20. The intensity is used as a free parameter in the model. The theoretical patterns for a fitting intensity of 80 mW/cm^2 are shown in Figure 4.7 and they agree well with the experimental diffraction patterns shown in Figure 4.6. We evaluate the averaged intensity of the incoming beam $I_{avg} = \frac{P}{\pi w_0^2}$, where w_0 is the waist of the beam and is measured to be 0.6 mm and find that the averaged intensity over the whole beam area is about 54 mW/cm^2 . The peak intensity of the incoming beam is also given by $I_{peak} = \frac{2P}{\pi w_0^2} \simeq 108 \text{ mW/cm}^2$ and the fitting intensity is right in between the averaged intensity and the peak intensity. When the standing wave is turned on, the atoms are travelling through the center of the beam where the intensity is higher compared to the edge of the beam but less than the intensity at the center; therefore, the atoms feel an

Interaction time

$t=0 \mu\text{s}$



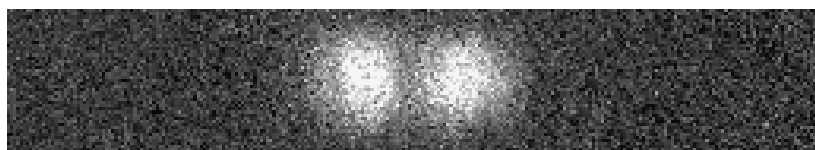
$t=2.16 \mu\text{s}$



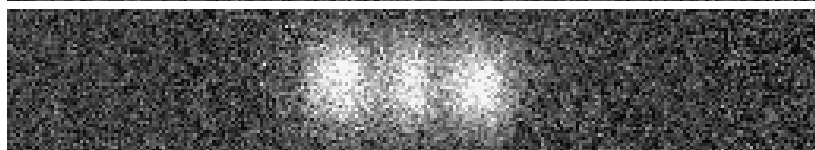
$t=4.12 \mu\text{s}$



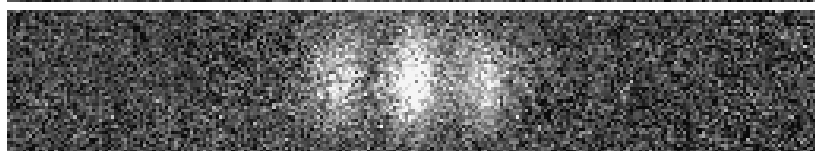
$t=6.04 \mu\text{s}$



$t=7.88 \mu\text{s}$



$t=9.92 \mu\text{s}$



$\longleftrightarrow$
185 μm

Figure 4.6: Absorption images of various interaction times. At zero interaction time, the condensate starts at zero momentum. When the interaction time increases, atoms start to populate higher momentum states, which are spatially separated after 8 ms propagation.

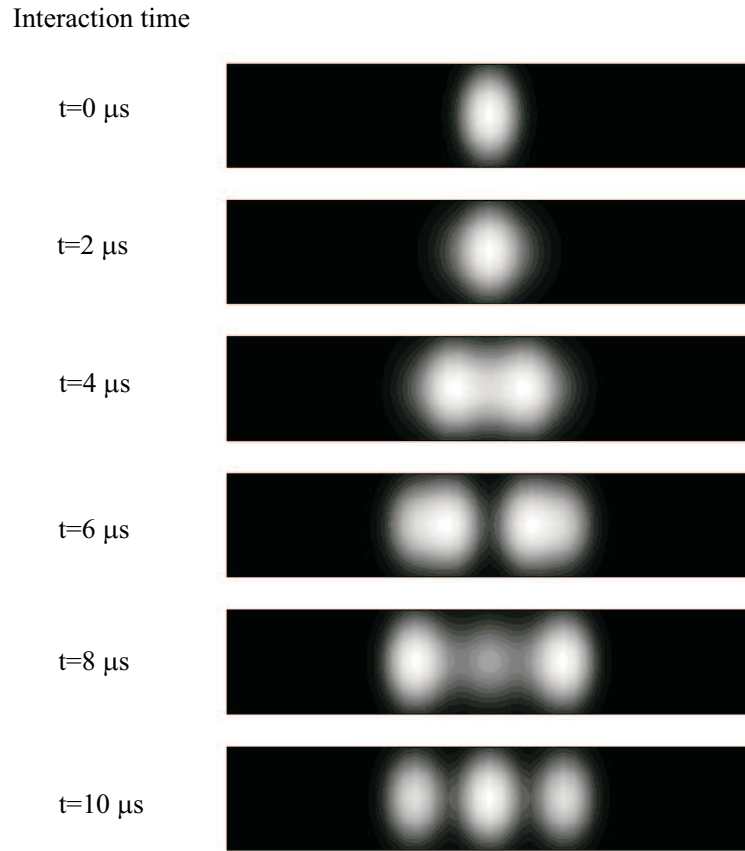


Figure 4.7: Theoretical diffraction patterns of various interaction times. The cloud sizes for the different diffraction orders are assumed the same as the size of the initial cloud. The intensity of the diffracting beam is a free parameter and is chosen to have a value of 80 mW/cm^2 to generate the results.

intensity which is higher than the averaged value but less than the peak value.

Because the diffraction pattern depends on both the interaction time and the intensity of the beam, because of the effective Rabi frequency $\Omega_{eff} \propto \Omega_0^2 \propto E_0^2 \propto I_0$, we could also vary the power ($P \propto I$) of the diffracting beam to affect the diffraction. Figures 4.8 shows the diffraction patterns for different powers of the diffraction beam and those diffraction patterns are similar to the patterns observed as we vary the interaction time. The interaction time is $4.16 \mu s$ and the fractions in the higher orders increase as we increase the power of the diffracting beam. The maximum diffraction order observed is the fourth order in Figure 4.8.

From the above experiment results, the diffraction pattern and efficiency qualitatively agree with the theoretical model. In order to achieve a quantitative comparison and also develop a guiding atom interferometer, in the following section we show this diffraction in the waveguide.

4.3 Diffraction of atoms in a waveguide

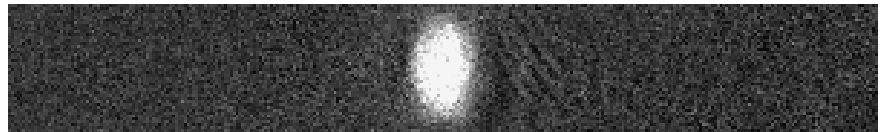
One of the main goals of this thesis is the development of a guided atom interferometer, the key element of which is a coherent beamsplitter for guided atoms. As I will show in this chapter, the optical beamsplitter described so far does indeed work for guided atoms. In addition, this also allows us to make a quantitative comparison with theory, because the guided atoms do not fall due to gravity, so that images can be taken with larger separation of the clouds.

4.3.1 A splitting pulse in the Raman-Nath regime

We perform the diffraction experiment in the waveguide with the setup described in Chapter 2. After a BEC is formed in the microtrap on the chip, we reduce the longitudinal confinement and release the BEC into the magnetic waveguide. The diffracting beam close to the surface is then turned on to split the cloud. After 38-ms propagation

Power of the
diffracting beam

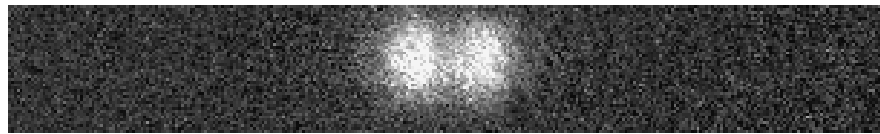
P=0.00 mW



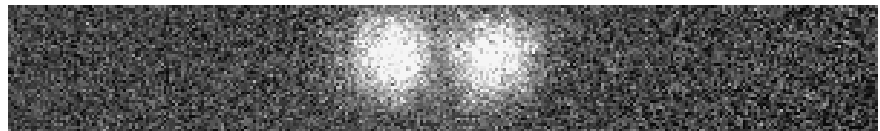
P=0.43 mW



P=0.98 mW



P=1.53 mW



P=1.80 mW

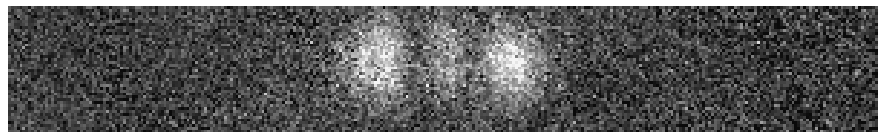


Figure 4.8: Diffraction patterns of different powers of the diffraction beam. The interaction time is $4.16 \mu\text{s}$.

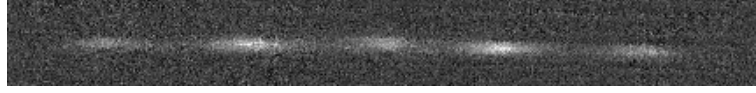


Figure 4.9: A diffraction pattern in the waveguide. The image is taken after 38 ms propagation.

in the waveguide, the different diffraction orders are spatially separated as shown in Figure 4.9. The center component is the zeroth order and we observe the symmetric higher orders up to the second order diffraction. Because the clouds are now spatially separated, we can fit each component with a Gaussian profile and evaluate the number in each diffraction order.

We again vary the interaction time, take a series of images, and then extract the numbers of atoms in each diffraction order from the images. The efficiency can be calculated by dividing the number of atoms in each component by the total number. We plot the number in each order for various times in Figure 4.10. The experimental data is fit to Bessel functions with the effective Rabi frequency as a variable. We plot the theoretical curve for different diffraction orders as the solid line shown in the Figure. In general, the data agrees very well with the prediction with a deviation of less than 20% diffraction efficiency from the theoretical curves. For a longer interaction time, the energy uncertainty associated with the pulse width becomes comparable to the energy splitting between the different momentum states and the assumption for a Raman-Nath pulse is not valid. As a result, we see the experimental data for longer interaction times in Figure 4.10 deviates from the theoretical prediction.

4.3.2 A reflecting pulse in the Bragg regime

The Raman-Nath pulse can split a single condensate into many wave packets through a coherent photon-scattering process. In order to make an interferometer, we need another pulse serving as a mirror to reflect the wave packets so that the wave

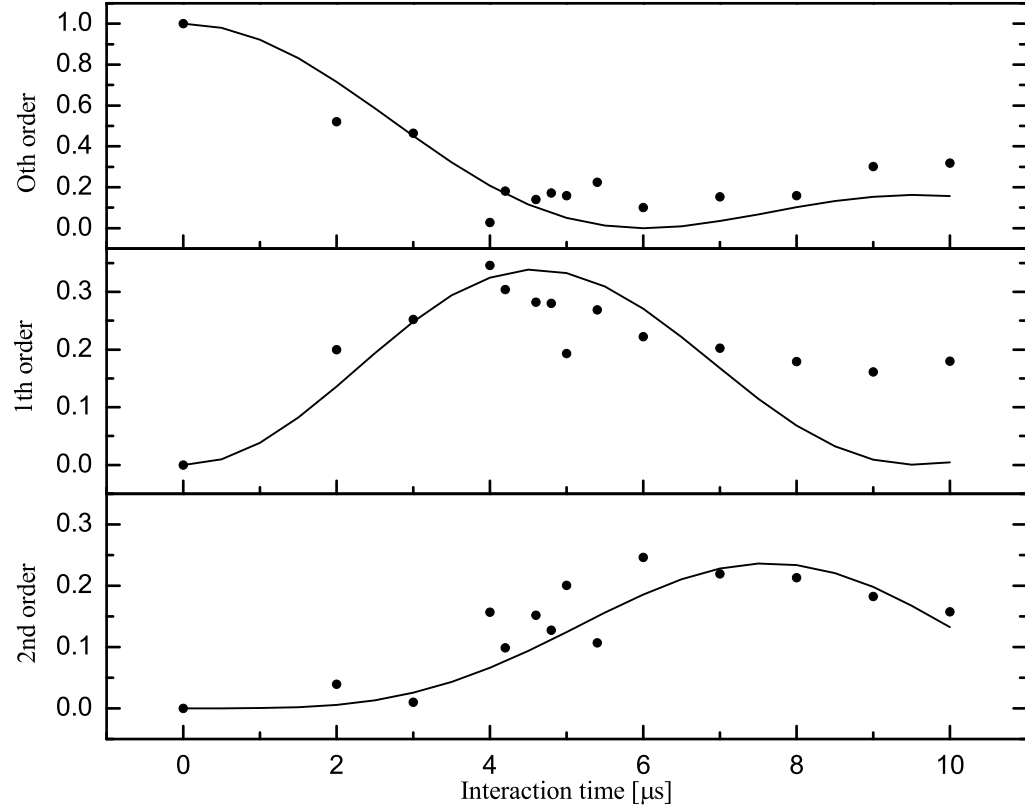


Figure 4.10: Diffraction efficiencies of various interaction times in the waveguide. The efficiencies for up to the second order are plotted. Solid circles are the experimental data and the lines are the fitting curves with an effective Rabi frequency of 400 kHz.

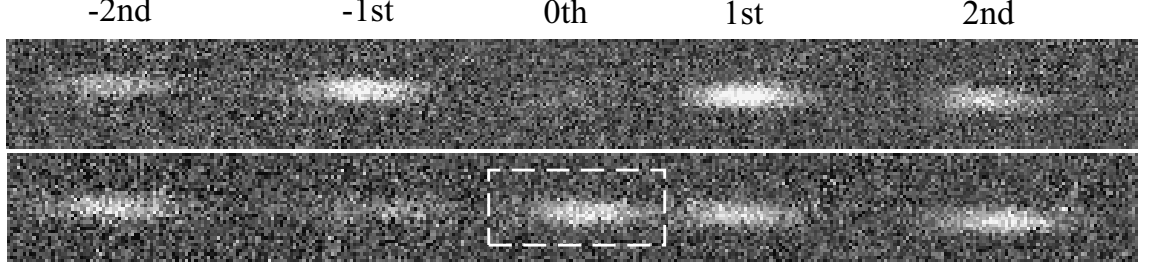


Figure 4.11: Diffraction patterns with and without a reflecting pulse. The images are taken 56 ms after the splitting pulse. The top image shows the diffraction pattern without applying a reflecting pulse. The bottom image is the diffraction pattern after applying a reflecting pulse. The reflecting pulse is turned on in between the splitting pulse and the imaging pulse with equal temporal separations, i.e., 28 ms after the first splitting pulse. The dashed rectangle marks the atoms that are reflected.

packets can recombine and interfere. This reflecting pulse can be done with a relatively long pulse in the Bragg scattering regime. A Bragg pulse simply reverses the momentum of the clouds as described in Section 4.1.2. Bragg scattering is a velocity-sensitive process because the atoms with different velocities, i.e., in different momentum states, have different effective Rabi frequencies as described in Equation 4.28. We can use this velocity-dependent process to reflect only a part of the split cloud, e.g., only the first order.

We optimize the intensity and the pulse width of the splitting pulse to distribute atoms in the first and second output orders but not the zeroth order. The first and second order travel outward with a momentum of $\pm 2\hbar k$ and $\pm 4\hbar k$, respectively. We purposely keep the longitudinal confinement on and wait for the atom velocities to slow down to half of their initial value, i.e., $\pm \hbar k$, before applying the Bragg pulse. This Bragg diffraction is a two-photon scattering process that changes the atom momentum from $\pm \hbar k$ to $\mp \hbar k$. The timing of the Bragg pulse can be estimated by the longitudinal oscillation period as $t_{Bragg} = \frac{T_{period}}{6} \sim 28$ ms, where T_{period} is the oscillation period of the harmonic trap. After optimizing the intensity and the pulse width of the reflecting pulse, the experimental result is shown in Figure 4.11. The images are taken 56 ms

after the splitting pulse and the Bragg pulse is applied at the half time between the splitting and the imaging pulses, i.e., 28 ms after the splitting pulse. The top image shows the diffraction pattern without applying a Bragg pulse and the bottom image shows the diffraction pattern after applying a Bragg pulse. The atoms shown at the center of the bottom image are the atoms reflected by the Bragg pulse. Theoretically, the Bragg scattering process could have a 100% coupling efficiency and for a perfect coupling efficiency, no atoms is left in the two first orders, which are the two components right next to the center component. However, a significant number of atoms are still observable in the first orders after fully optimizing the Bragg pulse. We also see the asymmetric distribution in the two first orders with the right component being stronger than the left one.

In order to find the origin of this asymmetry, we measure the velocity of each component at the time when the Bragg pulse is turned on. The velocities for the two first orders are 0.26 (to the right) and -0.58 (to the left) cm/s and the velocities for the two second orders are 0.75 and -1.04 cm/s. Those velocities are very different from what we expected. The momentum of the first and the second order are expected to be $\hbar k = 0.59$ and $2\hbar k = 1.18$ cm/s, so it seems that atoms going to the right side are slower than the atoms going to the left. We measure the initial velocity of the condensate before applying any pulse and the initial velocity is observed to have a maximum value of about 0.5 cm/s. When the cloud has an initial velocity (v_0), this velocity is added to the velocity transferred from the photons. If we consider only the first order diffraction, the atoms going to the right (or left) will have a net momentum of $m_{Rb}v_0 + 2\hbar k$ (or $m_{Rb}v_0 - 2\hbar k$), respectively. The atoms propagating in the opposite direction do not have the same speeds.

The initial velocity is due to a slosh motion of the center of the condensate in the trap after the condensate is created. This occurs because we need to move atoms farther away from the surface so the atom cloud is centered at the standing wave. The center

of the condensate is shifted by ramping the current in the wires and the transverse bias. A slosh motion is produced if the ramp is too fast. We can reduce this slosh motion by ramping the current more slowly or making a condensate far from the surface. After reducing the initial velocity, the asymmetry of the transfer efficiencies is also reduced.

In addition to the asymmetry in the transfer efficiencies, we also see that the split clouds slow down more than we expect as they climb up the potential. We attribute this to the anharmonicity of our longitudinal confinement. In our case, the ideal velocities to have the best Bragg scattering efficiencies are $p = \pm \hbar k$, and we optimize the velocities of the clouds by changing the timing of the Bragg pulse from 28 ms to 25 ms.

After those adjustments, the velocity of the two first orders are 0.54 and -0.57 cm/s, and the velocity of the two second orders are 1.13 and -1.07 cm/s.

If we keep the pulse width as a constant, e.g., 196 μs , and vary the pulse intensity, the probability of atoms in the zeroth and first orders will vary in a sinusoidal way. The fractional numbers in different orders vs various intensities are plotted in Figure 4.12. The oscillation, which can be observed between the zeroth and first order state up to 3 cycles, is called the Rabi Oscillation, or “Pendellösung” Oscillation [47, 48, 49, 50, 51], and the number in the second order remains constant as the intensity increases. We can optimize the coupling efficiency by varying the length and the power of the reflecting pulse. A nearly 100% coupling efficiency is achieved by a reflecting pulse of 150 μs in duration and around 6.2 μW in power. This optimized Bragg pulse is used to reflect the split cloud in an interferometer as discussed in Chapter 6.

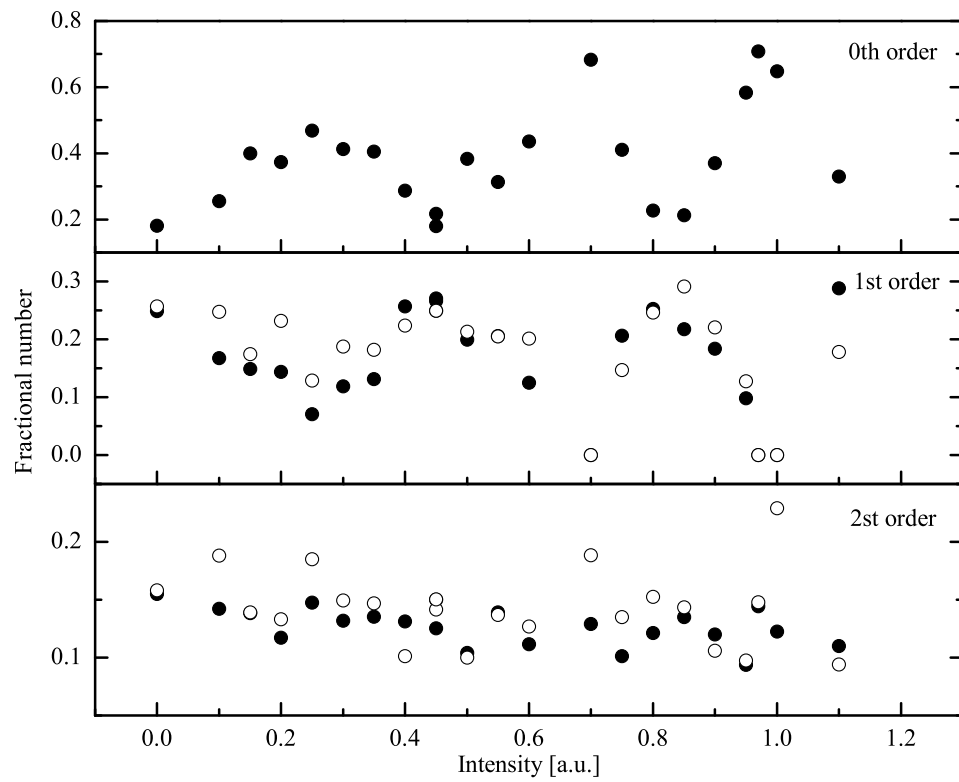


Figure 4.12: Efficiencies of various intensities for a first-order Bragg-diffraction. The Pendellösung oscillation is visible for up to 3 Cycles.

Chapter 5

An ideal beamsplitter

The diffracting pulse introduced in the previous chapter coherently splits the atom cloud into many wave packets. The diffraction efficiency to the m_{th} momentum state is given by a Bessel function of the m_{th} order. If we plot the diffraction efficiencies for the first three orders, i.e., 0th-2nd order, as functions of the intensity, we find that more than two wave packets exist for any non-zero intensity. The first diffraction order has the best efficiency of 34% at $\Omega_{eff}t = 1.84$ to observe the interference. However, the presence of atoms in higher or zero momentum states can reduce the visibility of the interference, if those atoms are partially reflected and recombined at the origin.

For an ideal beamsplitter, the condensate is equally split into two wave packets that propagate in opposite directions. A bichromatic standing wave that can reach an ideal 50/50 splitting ratio has been used in free space experiments [7, 8]. In this chapter, I will first discuss the application of such a bichromatic beamsplitter on our atom chip, and then introduce a new scheme developed in this thesis, namely a double-pulse beamsplitter. While it turns out that the bichromatic scheme is problematic in our setup, we have been able to successfully employ the double pulse scheme in the experiment to reach a 50/50 splitting ratio.

5.1 A double-frequency beamsplitter

Most free-space experiments use a double-frequency Bragg pulse to split BECs [7, 8]. The Bragg pulse can be made as a $\frac{\pi}{2}$ pulse so the BECs are equally split into two wave packets, which have different momenta, and later recombined to form an interferometer. The frequencies of the two counterpropagating beams that make the light waves are detuned from the atomic resonance. A frequency difference is introduced equal to $4\frac{\omega_{rec}}{2\pi} = 4\nu_{rec}$ between the two beams. These two counterpropagating waves are equivalent to a travelling wave with a beat frequency of $4\nu_{rec}$, which is equal to the energy spacing between the $|p = 0\rangle$ and $|p = 2\hbar k\rangle$ states. Both the energy and momentum of the atom-photon system are conserved in the splitting process. As a result, atoms equally populate the $|p = 0\rangle$ and $|p = 2\hbar k\rangle$ states. This splitting is not symmetric in terms of the momenta of the clouds, i.e., atoms only populate the $|p = 2\hbar k\rangle$ but not the $|p = -2\hbar k\rangle$ state, because of the asymmetry of the bichromatic standing wave configuration, as shown in Figure 5.1(a).

On our atom chips, the standing light wave is created by a laser beam retroreflected by a mirror. The frequencies of the incident and the reflected beam can not be controlled separately. We can modulate the amplitude of the rf signal coupled into the acousto-optic modulator, which is used as a switch for the standing light wave, with a frequency of $\frac{\Delta f}{2}$ to make a bichromatic standing light wave. This amplitude modulation creates two sidebands in the spectrum of the standing light wave. The frequencies of two sidebands are $f \pm \frac{\Delta f}{2}$, where $f = \omega/2\pi$ is the frequency of the laser beam without the modulation. If we fully modulate the amplitude of the rf signal coupled into the acousto-optic modulator, the output frequencies will have only the two frequency components at $f \pm \frac{\Delta f}{2}$. Therefore, we create a symmetric bichromatic standing light wave with a frequency difference of Δf , as shown in Figure 5.1.

The idea to coherently split atoms based on a bichromatic standing light wave was

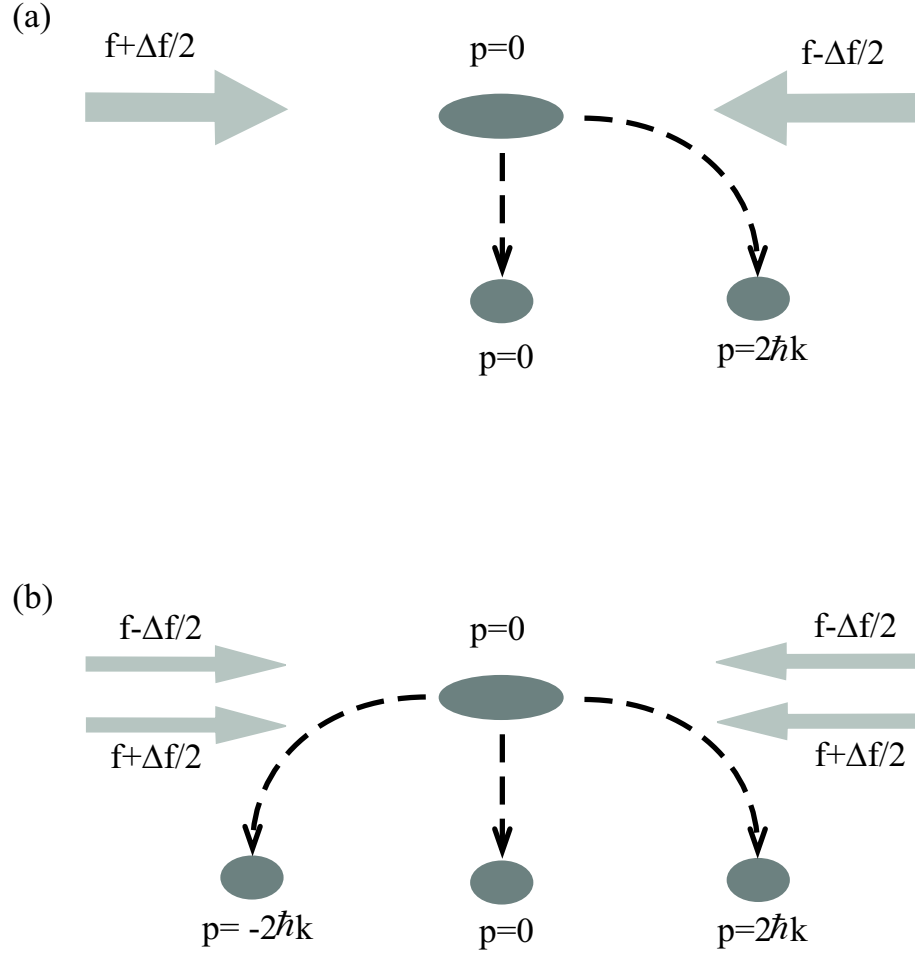


Figure 5.1: Bichromatic standing light waves. Bichromatic standing light waves in (a) an asymmetric configuration and (b) symmetric configuration. The frequency difference between the two components is $4\nu_{rec}$.

first proposed by Grimm et al. [52]. They suggested creating a ‘blazed’ phase grating beamsplitter rather than a sinusoidal phase grating. Analogous to an optical grating, a blazed phase grating has a higher diffraction efficiency than one with a sinusoidal profile. The function of the bichromatic standing wave that has a spatial phase shift between the two frequency components can be understood by considering the equivalent situation of two counterpropagating-travelling waves with time-dependent intensities. Atoms see the beam from one direction and then from the opposite direction after a time delay. The best diffraction efficiency occurs when the atoms experience a series of counterpropagating π pulses [53]. In the first pulse, atoms absorb a photon and receive a momentum of $\hbar k$ in one direction. The next pulse comes after $t = \frac{\pi}{\Delta}$ from the opposite direction. Atoms are stimulated to emit another photon in the opposite direction; therefore, atoms gain a net $2\hbar k$ momentum after each cycle. For first-order diffraction with a momentum of $2\hbar k$ transferred from the photons to the atoms, the interaction time should be $t = \frac{2\pi}{\Delta}$, which is equal to the duration of one cycle.

The case of a bichromatic standing wave was analyzed in detail by Tan et al., [53] and Johnson et al. [54] who found that the maximum efficiency of splitting can only be achieved when the spatial phase shift of the two standing waves equal $\pi/2$. However, we cannot reach this spatial phase shift between the two bichromatic standing waves on our atom chip. Our standing wave is created by the retroreflection from one of the prism mirrors, and the electric field has a node at the location of the mirror. The distance (d) from the location of the cloud to the mirror is about 2 mm. To create a $\frac{\pi}{2}$ spatial phase shift at the location of the condensate, a frequency difference of $\Delta = 2\pi \frac{c}{4d} = 2\pi \times 37.5$ GHz is required, where c is the speed of the light. This required detuning is beyond our tuning range.

5.2 A double-pulse beamsplitter

To achieve an ideal beamsplitter, we develop a monochromatic double-pulse scheme, which involves a relatively simple setup. The details of the double-pulse beamsplitter are discussed in the following section. The splitting pulse consists of two subpulses with the same frequency and each pulse works in the Raman-Nath regime. To evaluate the required pulse width and the intensity, we use the coupled equations described in Chapter 4.

If the initial velocity of the cloud is zero [55], the probability of atoms in each momentum state can be described as

$$i\hbar \frac{\partial g_m(t)}{\partial t} = (m^2 \hbar \omega_{rec} + \hbar \Omega_{eff}) g_m(t) + \frac{\hbar \Omega_{eff}}{2} [g_{m+2}(t) + g_{m-2}(t)], \quad (5.1)$$

where $g_m(t)$ is the probability of the atoms in the m_{th} momentum state. If we choose a weak intensity ($\frac{\hbar \Omega_{eff}}{2} \ll 16 \hbar \omega_{rec}$), we can ignore the probability of finding atoms in the fourth or higher momentum states, i.e., $g_m \simeq 0$ when $m \geq 4$.

We only consider the coupled equations up to the first order. If we ignore the constant phase evolution term, $\hbar \Omega_{eff} g_m(t)$, for all the different momentum states, those coupled equations can be rewritten as

$$\begin{aligned} i \frac{\partial g_0(t)}{\partial t} &= \frac{\Omega_{eff}}{2} [g_2(t) + g_{-2}(t)] \\ i \frac{\partial g_2(t)}{\partial t} &= 4\omega_{rec} g_2(t) + \frac{\Omega_{eff}}{2} g_0(t) \\ i \frac{\partial g_{-2}(t)}{\partial t} &= 4\omega_{rec} g_{-2}(t) + \frac{\Omega_{eff}}{2} g_0(t) \end{aligned} \quad (5.2)$$

Based on symmetry, we can define $g_+(t) = \frac{1}{\sqrt{2}} [g_2(t) + g_{-2}(t)]$ (symmetric term) and $g_-(t) = \frac{1}{\sqrt{2}} [g_2(t) - g_{-2}(t)]$ (asymmetric term). The coupled equations become

$$\begin{aligned} i \frac{\partial g_0(t)}{\partial t} &= \frac{\Omega_{eff}}{\sqrt{2}} g_+(t) \\ i \frac{\partial g_+(t)}{\partial t} &= 4\omega_{rec} g_+(t) + \frac{\hbar \Omega_{eff}}{\sqrt{2}} g_0(t) \\ i \frac{\partial g_-(t)}{\partial t} &= 4\omega_{rec} g_-(t). \end{aligned} \quad (5.3)$$

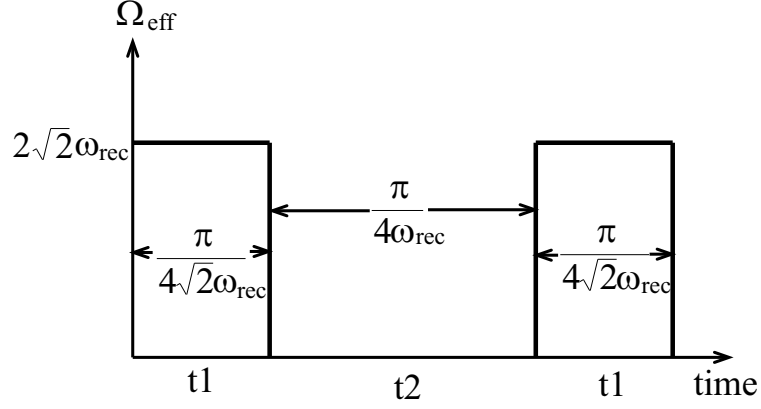


Figure 5.2: The sequence of a double-pulse scheme. The splitting pulse consists of two subpulses with the Rabi frequency $\Omega_{eff} = 2\sqrt{2}\omega_{rec}$, the pulse width $\Delta t_1 = \frac{\pi}{4\sqrt{2}\omega_{rec}}$, and the time interval between the two subpulses $\Delta t_2 = \frac{\pi}{4\omega_{rec}}$.

If all the atoms are initially at $|p = 0\rangle$, i.e., $g_{\pm}(0) = 0$, $g_{-}(t)$ is zero for any time t , the last equation vanishes and the coupled Equations are identical with a two-level system.

5.2.1 Bloch diagram representation

This two-level coupling system can be described with a graphical representation called a ‘Bloch diagram’ [56]. In a Bloch diagram, we use a Rabi vector to represent the coupling strength, i.e., the Rabi frequency Ω_{eff} , and a state vector to describe how the quantum states of the atoms evolve. We can get a 100% transfer efficiency from the $|p = 0\rangle$ to $|p = \pm 2\hbar k\rangle$ state by applying two identical pulses with $\Omega_{eff} = 2\sqrt{2}\omega_{rec}$, pulse width $\Delta t_1 = \frac{\pi}{4\sqrt{2}\omega_{rec}}$, and the time interval between the two pulses $\Delta t_2 = \frac{\pi}{4\omega_{rec}}$, as shown in Figure 5.2. On the Bloch sphere, the two states are on the same axis (see Figure 5.3) with the north pole representing the first momentum state, $|p = \pm 2\hbar k\rangle$, and the south pole representing the zero momentum state, $|p = 0\rangle$. The energy difference between the zero momentum and the first momentum state is $4\hbar\omega_{rec}$. Atoms start initially at $|p = 0\rangle$, so the state vector points down. When the standing wave is turned

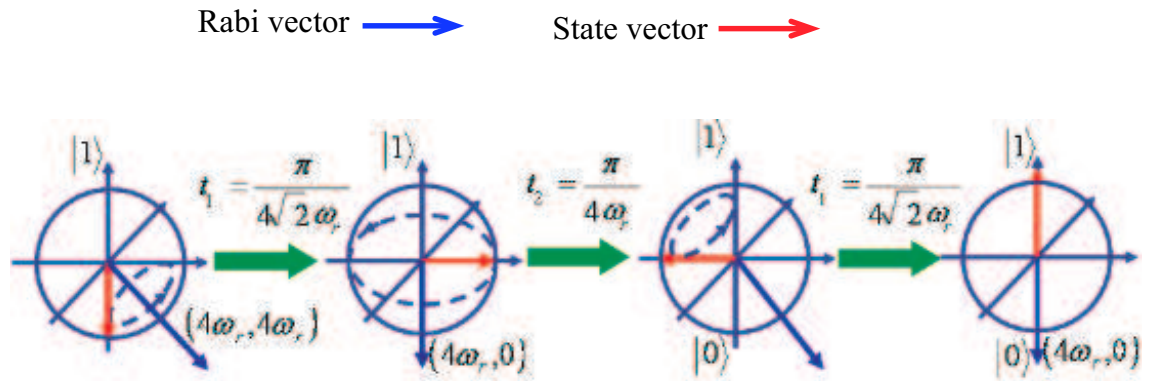


Figure 5.3: Bloch diagram for a double-pulse beamsplitter. When the first pulse is turned on, the Rabi vector points 45° and the state vector precesses to a perpendicular position. After the first pulse, the state vector accumulates a π phase shift until the second pulse is turned on. During the second pulse, the state vector precesses again along the 45° -oriented Rabi vector and points up when the second pulse is turned off.

on with the Rabi frequency equal to $2\sqrt{2}\omega_{rec}$, the Rabi vector points 45° with respect to the initial state vector. The state vector then precesses along the Rabi vector. After time $t_1 = \frac{\pi}{4\sqrt{2}\omega_{rec}}$, the state vector becomes perpendicular to the axis, and the atoms are then in a superposition of the zero and first momentum state, i.e., atoms have equal probabilities of being in the zero and first momentum states. After the first pulse is turned off, the phase of the atoms evolve and the second pulse is turned on when a π phase is accumulated, which is after $t_2 = \frac{\pi}{4\omega_{rec}}$. During the second pulse, the state vector precesses again along the 45° -oriented Rabi vector and more atoms are coupled into the first momentum state. After the second pulse is turned off, the state vector points up, which means that all the atoms are now in the first momentum state. The optimum condition for the two pulses is not unique but using two identical pulses is the simplest way to obtain a maximum splitting efficiency. The effective Rabi frequency in this case satisfies the weak intensity condition ($\sqrt{2}\omega_{rec} \ll 16\omega_{rec}$), so the populations in the higher momentum states are negligible. This double-pulse scheme is also expected to have a very high efficiency even for coupling the atoms to a higher momentum state and a more detailed discussion is given in Ref. [55].

5.2.2 Experimental results

In our experiment, we obtain an optimized splitting ratio of nearly 100% within our detection limit as shown in Figure 5.4. No visible atoms remain in the zero momentum state after the double-pulse sequence.

We can also vary the delay between the two pulses to change the population ratio of the two states. If we let a 2π instead of a π phase accumulate between the two pulses, the final state vector points down instead of pointing up, because all the atoms go back to the zero momentum state after the double pulses. If a phase between 0 and π has accumulated, the final state vector will be off axis causing the atoms to be distributed in the $|p = 0\rangle$ and $|p = \pm 2\hbar k\rangle$ states. We expect the atom population will move from

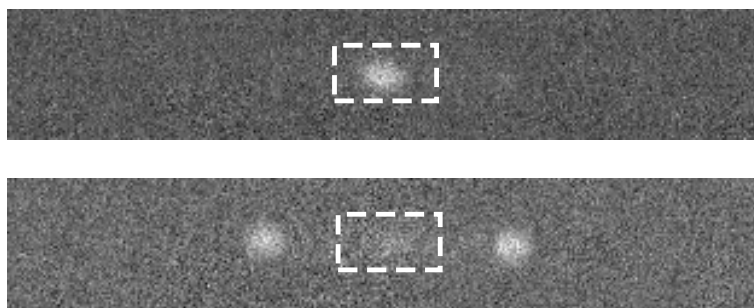


Figure 5.4: Absorption images of an ideal Beamsplitter. Upper (lower) image shows the absorption image before (after) applying the double pulse. The dashed line marks atoms in the zero momentum state. After the double pulse, atoms are equally distributed in the plus and minus first orders.

one momentum state to the other and comes back as we increase the delay between the two pulses. The oscillation period is equal to $T = \frac{2\pi}{4\omega_{rec}} \sim 66 \mu\text{s}$. The experimental results are shown in Figure 5.5. The populations in the $|p = 0\rangle$ and $|p = \pm 2\hbar k\rangle$ states are anticorrelated with atoms equally distributed in both the $|p = 2\hbar k\rangle$ and $|p = -2\hbar k\rangle$ states.

5.2.3 Velocity-dependent splitting ratios

Up to now, I have considered the case where the cloud is initially at rest. The situation changes when the cloud has an initial velocity. The splitting ratio is affected by the initial velocity. We consider a cloud with an initial velocity of $v_{in} = \frac{\hbar k_{in}}{m_{Rb}}$ and the coupled equations then become

$$\begin{aligned} i\frac{\partial g_0(t)}{\partial t} &= \frac{\Omega_{eff}}{\sqrt{2}}g_+(t) \\ i\frac{\partial g_+(t)}{\partial t} &= 4\omega_{rec}g_+(t) + \frac{\hbar\Omega_{eff}}{\sqrt{2}}g_0(t) + \frac{4k_{in}}{k}\omega_{rec}g_-(t) \\ i\frac{\partial g_-(t)}{\partial t} &= 4\omega_{rec}g_-(t) + \frac{4k_{in}}{k}\omega_{rec}g_+(t), \end{aligned} \quad (5.4)$$

where k is the wave vector of the photons. The asymmetric term, $g_-(t)$, is no longer decoupled from the other two equations. As we increase the initial velocity, the coupling between symmetric and asymmetric states becomes stronger; therefore at the same time, the splitting ratio with an initial condition of $g_0(0) = 1$ and $g_{\pm}(0) = 0$ becomes more asymmetric, as shown in Figure 5.6. To achieve an ideal beamsplitter, the initial motion of the cloud must be minimized. The initial motion can be efficiently reduced usually by slowly varying the potential when we release the condensate from the trap to the waveguide and a nearly ideal 50/50 beamsplitting ratio can be achieved. In the next chapter, I will describe an experiment where an initial velocity is required to observe the interference and the asymmetry of the beamsplitter needs to be reduced in order to have a greater contrast ratio. This asymmetry can be suppressed by using shorter pulses because the energy uncertainty arising from the finite pulse width is larger when

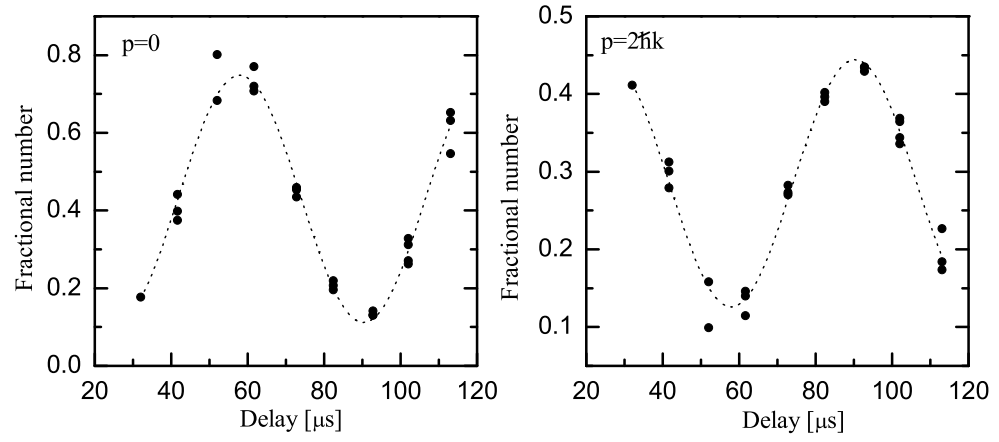


Figure 5.5: Populations of atoms in the $|p=0\rangle$ and $|p=2\hbar k\rangle$ states. The populations in the both states are anticorrelated because of number conservation.

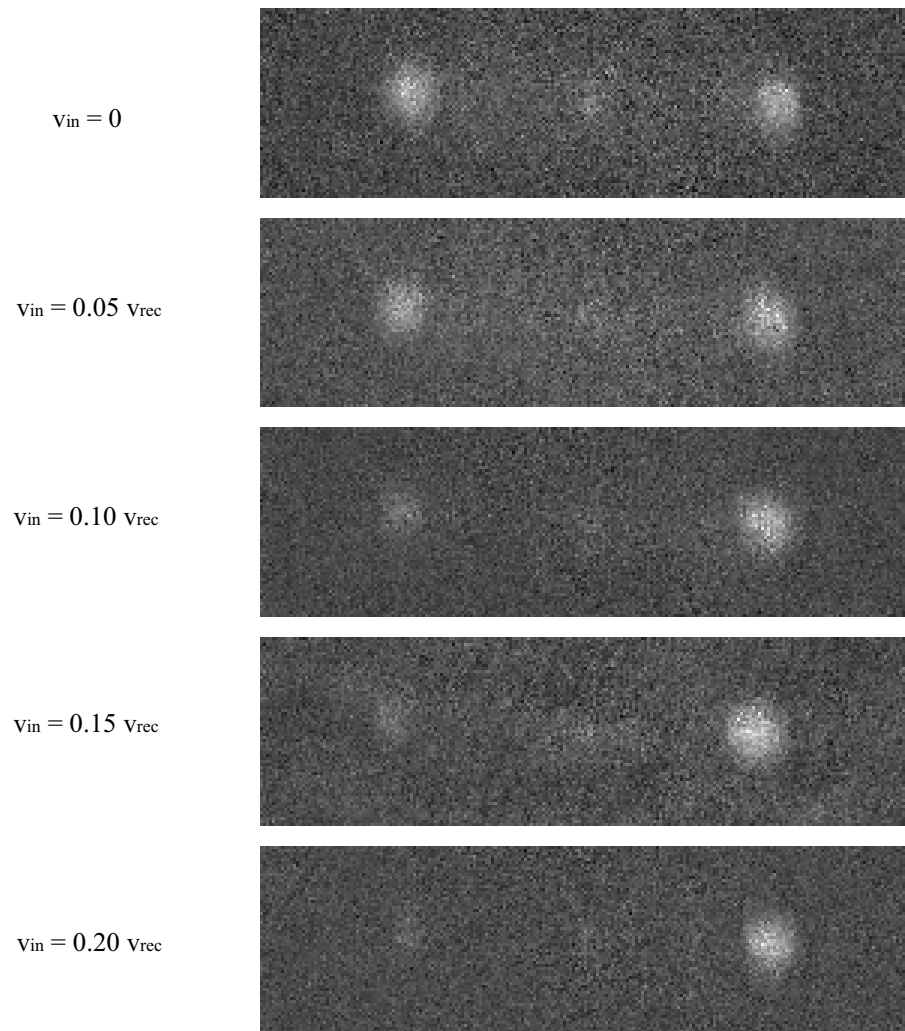


Figure 5.6: Asymmetric splitting ratios due to non-zero initial velocities. The asymmetry increases as we increase the initial velocities.

the pulse is shorter. Thus the splitting ratio is less sensitive to energy differences between the two momentum states caused by the initial motion. If the cloud has an initial velocity of about $0.2 v_{rec}$, when we change the pulse width from $50 \mu s$ to $10 \mu s$, the beamsplitting ratio changes from $10\backslash 90$ to $45\backslash 55$ after optimizing the intensity to minimize the number in the zero momentum state. In general, a shorter splitting pulse is preferred in the experiment so the splitting ratio can be less sensitive to the initial velocity.

Chapter 6

An atom interferometer on a chip

In the previous chapter, I have described the implementation of a new double pulse beamsplitter. Here I will show how to extend this scheme to a fully functioning atom interferometer.

An atom interferometer includes the process of splitting, reflecting, and recombining the wave packets. We use the double-pulse scheme to create high-efficiency, coherent splitting, which will be referred to as “a splitting pulse” in the following discussion. A Bragg pulse serves as a mirror to reflect atoms. After the atoms are recombined, the two wave packets interfere, and absorption images are taken to observe the interference.

6.1 Interference patterns

After the splitting pulse, atom clouds propagate in opposite directions, each with a momentum of $\pm 2\hbar k$. Upon their return, they interfere once they are spatially overlapped. However, the two travelling wave packets have a relative velocity of $\frac{4\hbar k}{m_{Rb}}$. Thus the interference of the two clouds has a spatial period of $\lambda/4$, where λ is the wavelength of the standing light wave (typically around 780 nm in our experiment). The resulting 200-nm spatial period is much smaller than our imaging resolution so interference is not visible in our images. Instead of directly imaging the interfered density profile of the cloud, we apply a pulse to recombine the clouds and read out the interference pattern. The recombining pulse converts the spatial interference pattern into a momentum distri-

bution. As I will describe below, the beamsplitting pulse can be used as a recombining pulse to achieve this goal.

The phase accumulated along the path can be expressed in terms of the classical action S along the path:

$$\phi = \frac{1}{\hbar} S. \quad (6.1)$$

The classical action is defined in terms of the Lagrangian, which is

$$L[x, \dot{x}] = \frac{1}{2} m \dot{x}^2 - V(x) \quad (6.2)$$

for a particle with mass m in a potential $V(x)$. The phase accumulated along each path then becomes [57, 58, 59]

$$\begin{aligned} \phi &= \frac{1}{\hbar} \int L[x(t), \dot{x}(t)] dt \\ &= \frac{1}{\hbar} \int \left[m v \frac{dx}{dt} - \left(V(x) + \frac{m v^2}{2} \right) \right] dt \\ &= \frac{1}{\hbar} \int (p dx - H dt). \end{aligned} \quad (6.3)$$

where H is the Hamiltonian for the classical motion of the particles. If the potential is symmetric and the condensate starts from rest, the clouds move apart and come back to their origin with the same speed. The two split clouds accumulate the same phase before they overlap again; therefore, when the recombining pulse puts them back to rest, they constructively interfere in the zero momentum state. We would see the clouds back at rest as its original position (see Figure 6.1). However, if a nonuniform perturbation in the splitting direction is added the clouds, e.g., a potential gradient is applied, the clouds no longer see the same potential. They accumulate different phase when the gradient is turned on. For the special case where the differential phase shift is equal to π , the two clouds destructively interfere in the zero momentum state after recombining. The density of atoms in the zero momentum state is nearly zero because of destructive interference and the atoms appear in the $|p = \pm 2\hbar k\rangle$. For any differential phase shift accumulated between 0 and π , population will be distributed in both momentum states.

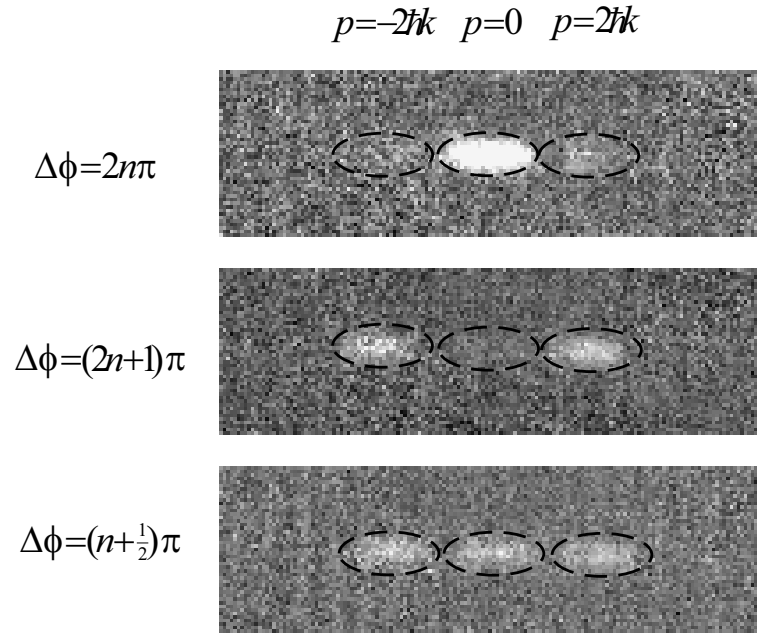


Figure 6.1: Interference patterns. The fractional numbers in the $|p=0\rangle$ and $|p=\pm 2\hbar k\rangle$ state are anticorrelated. The probability in each state depends on the differential phase shift.

Because of number conservation, the population in the each momentum state appears as an anticorrelation. The probability of atoms in each momentum state is given by

$$P_{p=0}(\phi) = \frac{1 + \cos(\phi)}{2} \quad (6.4)$$

$$P_{|p|=2\hbar k}(\phi) = \frac{1 - \cos(\phi)}{2}, \quad (6.5)$$

where ϕ is the phase difference accumulated between the split clouds.

6.2 A double-pulse interferometer

In order to demonstrate the read-out scheme, we first employ a simplified interferometer that consists of only two splitting pulses and no reflecting pulse. After the first pulse splits the condensate into two, the second pulse is turned on after a delay to recombine them. The delay between the two pulses has to be short enough for the two split clouds to remain spatially overlapped so they can interfere. Because the clouds are spatially overlapped, we can not apply a spatially dependent potential to introduce a phase shift. But we can introduce an initial velocity to create a difference in the kinetic energies. If the condensate has a zero initial velocity, the two split clouds have the same kinetic energy. They experience the same potential energy if the potential is symmetric. No phase difference is accumulated between the two clouds in this case. If the initial velocity of the condensate is not zero, after the first splitting pulse, one of the clouds has a higher speed than the other because the momentum is conserved during the splitting and the two clouds gains a $2\hbar k$ momentum but in opposite directions. Because the clouds have different momenta and kinetic energies, their phases vary according to Equation 6.3. The phase difference between the clouds propagating in opposite directions is proportional to the time delay. We observe this differential phase shift by increasing the delay between the splitting and recombining pulses. We introduce an initial velocity to the condensate by suddenly shifting the bottom of the trap right before releasing the condensate to the waveguide. This shift of the trap bottom

is done by applying a magnetic field gradient. Due to an asymmetry in the trapping field, this leads to an initial velocity of the cloud as described in [20]. The experimental results are shown in Figure 6.2. The initial velocity is $0.1v_{rec} = 0.59$ mm/sec, where $v_{rec} = \frac{\hbar k}{m_{Rb}}$ is the recoil velocity. The expected period of the sinusoidal oscillation is $675 \mu s$ as we increase the delay between the two pulses, and the measured period is around $660 \mu s$. We also change the period by increasing the initial velocity by a factor of two. The measured period becomes $280 \mu s$ compared to a expectation value of $337 \mu s$ (see Figure 6.3). The measured period is obtained by fitting the experimental data with a sinusoidal function.

We measure the initial velocity by observing the center displacement of the condensate as we increase the propagation time. Because the center displacement of the condensate is small compared to the cloud size, this measurement gives us a 10% error in our initial velocity measurement, which contributes to a 10% error in the estimation of the interference period.

The contrast ratio of the interference signals decays as we increase the delay. The reduction of the contrast ratio is mainly due to the spatial separation of the two clouds. After the splitting pulse, the clouds move apart and the spatial overlap of the two split clouds starts to decrease. As the delay increases, the overlap decreases. Once the clouds are spatially separated, they do not interfere and the contrast ratio decays to zero. The contrast ratio can be estimated by calculating the atom number in the overlap region. Assuming that the clouds have Gaussian density profiles, the contrast ratio, which is proportional the fractional number in the overlap area, is an error function of the center separation (x) between the clouds and is given by

$$\eta(x) = 1 - \text{erf}\left(\frac{x}{2\sigma_x}\right) = 1 - \text{erf}\left(\frac{4v_{rec}t}{2\sigma_x}\right), \quad (6.6)$$

where σ_x is the rms size of the cloud in the splitting direction. The typical rms size of the cloud is around $25 \mu m$ so we apply the modification factor, $\eta(x)$, to the sinusoidal

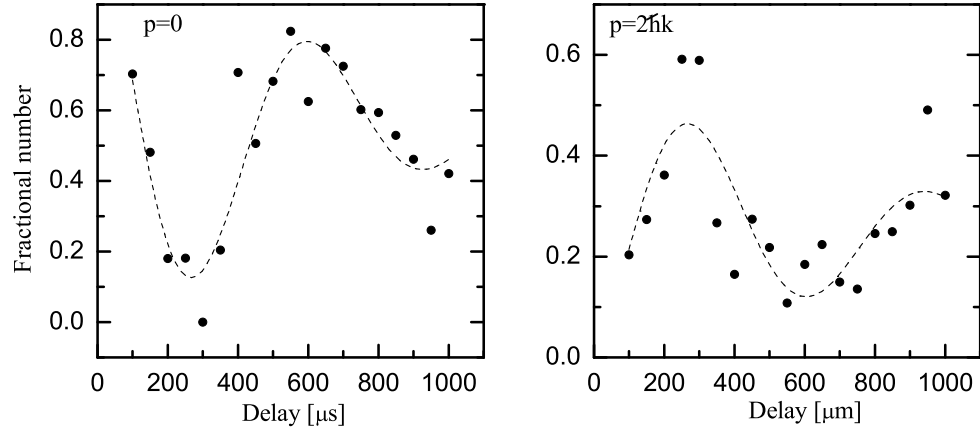


Figure 6.2: Interference signals of various delay times between the splitting and recombining pulses. The interference signals are plotted as the fractional number in the zero and the first momentum state. The initial velocity of the condensate is $0.1v_{rec}$.

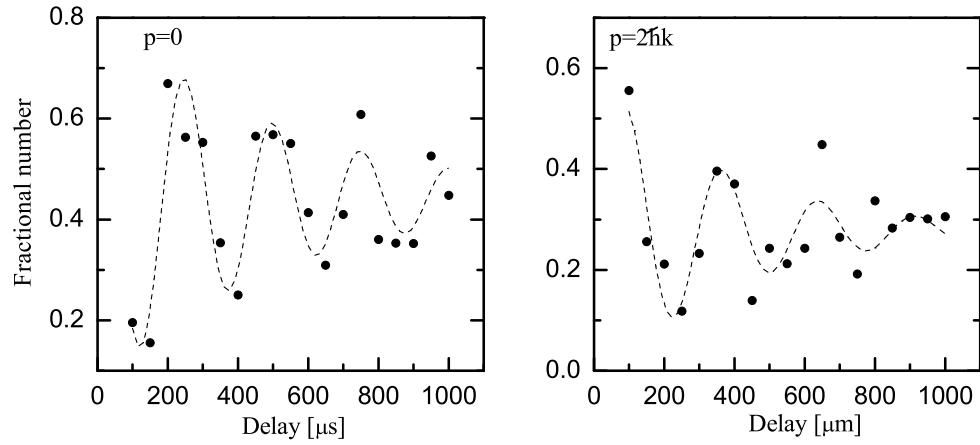


Figure 6.3: Interference signals with a different initial velocity of the condensate. The initial velocity of the condensate is about $0.2v_{rec}$.

interference curve. We compare the theoretical curve to the experimental data of the fraction of atoms in the zero momentum state in Figure 6.2, and obtain good agreement, as shown in Figure 6.4. The interference signal has been observed up to around 1 ms with the double-pulse interferometer, at which point the spatial separation has lead to a complete loss of contrast. This interferometer also proves that splitting and recombining are coherent processes, and the coherence of the condensate is preserved during the pulses. To observe interference for a longer propagation time, we need to include a reflecting pulse so that we can reflect the split clouds and make them come back to their initial position.

6.3 A Michelson interferometer

We develop an on-chip Michelson interferometer that employs a splitting pulse, a reflecting pulse, and a recombining pulse. In analogy to the optical Michelson interferometer, the splitting and the recombining occur at the same spatial location. Figure 6.5 shows the sequence of the Michelson interferometer. The clouds move apart after the splitting pulse, and the reflecting pulse reverses the momentum of the clouds. Upon their return to the original position, the recombining pulse is applied to observe the interference. We find the optimum splitting occurs when the two subpulses are both 20 μs in duration with a power of around 5.5 μW and the delay between the turning on time of the two pulses is 63 μs . The reflecting pulse is chosen to be 150 μs in duration with a power of 6.2 μW .

A differential phase shift has been introduced by two different ways: a magnetic gradient or an initial velocity of the condensate in a potential well. Experimental results are presented in the following sections.

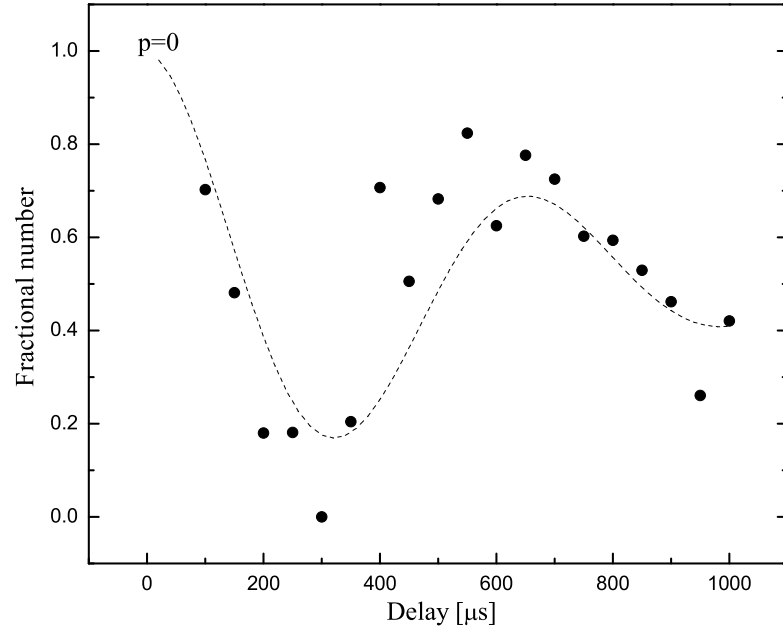


Figure 6.4: A comparison of the theoretical prediction and experimental measurement of interference signals. The solid circles are the fractional numbers of the atoms in the zero momentum state, as shown in Figure 6.2. The dashed line is the expected interference signal with the decay of the contrast ratio due to the spatial separation of the clouds.

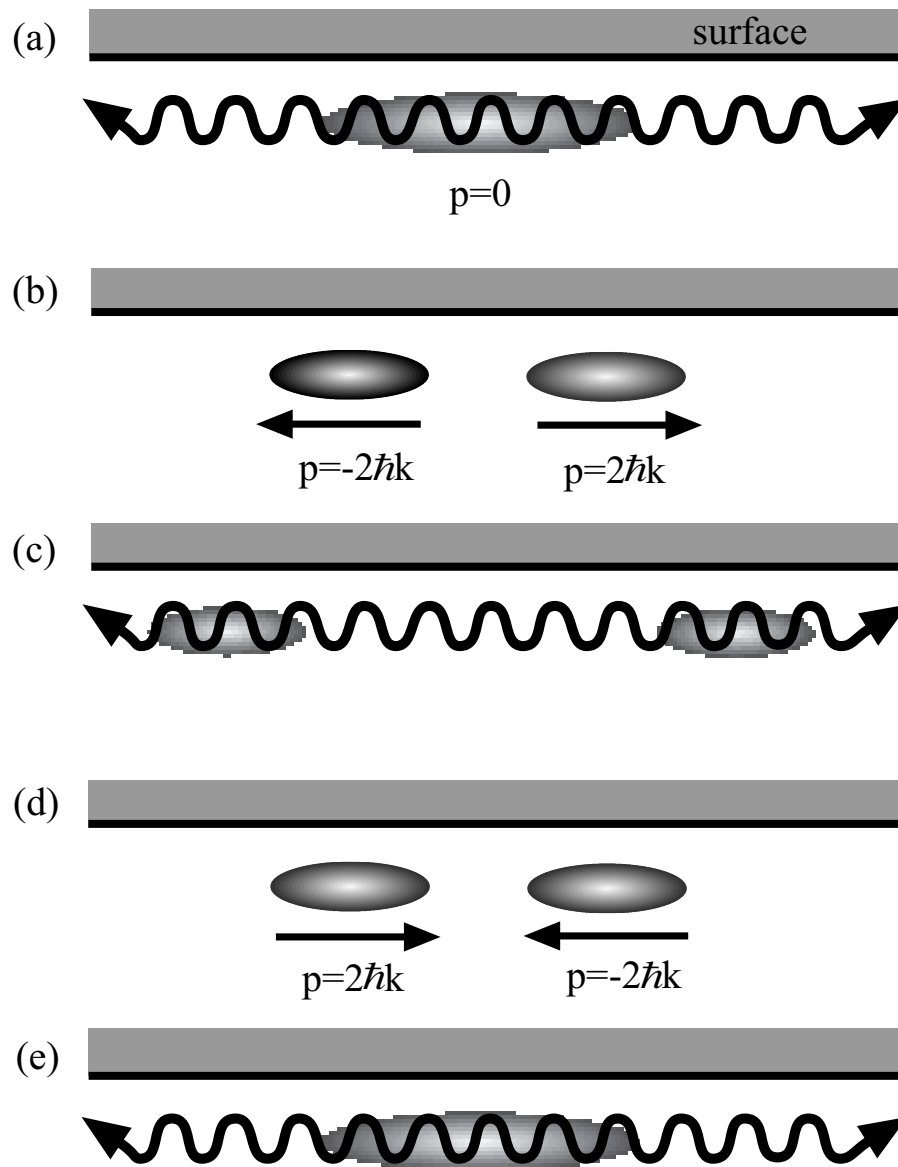


Figure 6.5: A Michelson interferometer. (a) The splitter pulse is turned on at $t=0$. (b) The condensate is split into two wave packets that propagate in the opposite direction at $0 < t < T/2$. (c) The reflection pulse is turned on at $t=T/2$. (d) The wave packets propagate back toward to the center of the waveguide at $T/2 < t < T$. (e) The recombining pulse is turned on at $t=T$ when the clouds are overlapped.

6.3.1 A phase shift introduced by a magnetic gradient

In the first experiment, we introduce a differential phase shift by a magnetic gradient. This magnetic gradient is provided by a current-carrying wire (called a “gradient wire”) that is perpendicular to the waveguide. The gradient wire sits 3 mm below the chip surface and is 2 mm offset horizontally from the initial location of the condensate. The total distance from the gradient wire to the condensate is 3.8 mm. The gradient wire produces a field $B(r) = \mu_0 \frac{I}{2\pi r}$, but since the cloud separation used in this experiment is small compared to the distance from the wire, the field can locally be approximated as a linear gradient B' . After the condensate is split by the first pulse, we apply a magnetic gradient pulse for a very short time. The differential phase shift is given by

$$\phi = \frac{1}{\hbar} \mu B' \int_{t_1}^{t_2} x(t) dt. \quad (6.7)$$

The differential phase is proportional to the separation of the clouds $x(t)$, magnetic gradient B' , and magnetic pulse width. For a larger separation, we get a better sensitivity of the differential phase shift. The interference of the condensates after 1 ms propagation time is shown in Figure 6.6. The maximum separation of the split clouds for a 1 ms total-propagation time in the waveguide is $12 \mu\text{m}$ which is the distance right before they are reflected. The splitting, reflecting, and recombining pulse are turned on at $t = 0, 0.5$, and 1.0 ms, respectively. The magnetic pulse is turned on $250 \mu\text{s}$ before the reflecting pulse for $500 \mu\text{s}$. Because the split clouds are moving when the magnetic pulse is turned on, the separation between the clouds is changing during the magnetic pulse. This average separation of the cloud during the magnetic pulse is $8.82 \mu\text{m}$. The period of the interference is expected to be 3.5 G/cm , which is close to the measured period of 3.8 G/cm within 10%. When we change the timing of the magnetic pulse, the average separation of the clouds during the magnetic pulse changes so that the period changes. Figure 6.7 shows the interference fringes after changing the timing of the

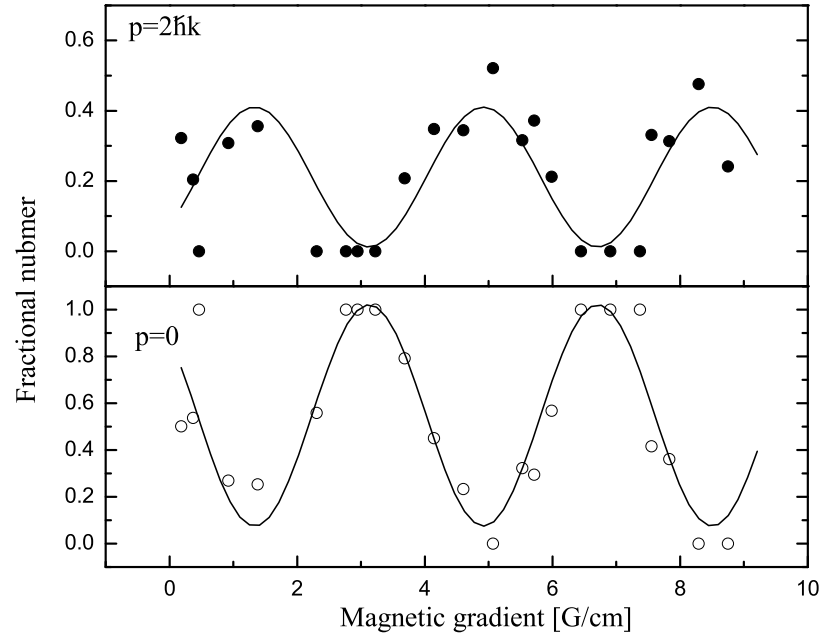


Figure 6.6: Interference after 1 ms propagation. The average separation of the clouds is $8.82 \mu\text{m}$ when the magnetic pulse is turned on. The expected period is 3.5 G/cm and the measured period is 3.8 G/cm.

magnetic pulse. The average separation of the clouds is estimated to be $6.94 \mu\text{m}$ and the expected period is 4.5 G/cm , which is also close to the measured period, 5.7 G/cm , with an error of 20%.

The measured periods are longer than the estimates because the actual separation is less than what we expect. The underestimation of the cloud separation is due to an unexpected residual harmonic trap. Ideally, the waveguide potential is homogenous in the guiding direction, but in practice, we find that a residual 5 Hz trap exists even after we turn down all the currents that provide the longitudinal confinement. This 5 Hz trap is likely caused by the imperfection of the guiding wire that creates a corrugated magnetic potential. The residual trap slows down the split clouds such that the separation between the clouds decreases. After accounting for the 5 Hz trap, the expected periods for the data in Figure 6.6 and Figure 6.7 become 3.7 G/cm and 5.3 G/cm , respectively, which agree well with the experimental results within less than 10% in both cases.

In addition to the residual trap, we observe an antisymmetric interference pattern in some of our experimental data. Figure 6.8 shows one example. In contrast to the expected symmetric density distribution in $|p = \pm 2\hbar k\rangle$ states, we see that a ‘second’ interference component appears in the $|p = \pm 2\hbar k\rangle$ states with a frequency of half of the normal interference frequency. This second interference component is due to an initial velocity of the cloud. When the cloud has a nonzero initial velocity, the splitting ratio is no longer 50/50 as described in Chapter 5, and a small number of atoms are also left in the zero momentum state after the splitting pulse. These leftover atoms are coupled into the $|p = \pm 2\hbar k\rangle$ states after the recombining pulse. They interfere with the atoms remaining in the $|p = \pm 2\hbar k\rangle$ after the three-pulse sequence causing the second interference component to occur. The leftover atoms stay at the origin when the magnetic pulse is turned on and the separation from the leftover atoms to the split clouds is only half the distance between the them. Therefore, the differential phase shift of this second interference effect is only half the normal phase difference between the

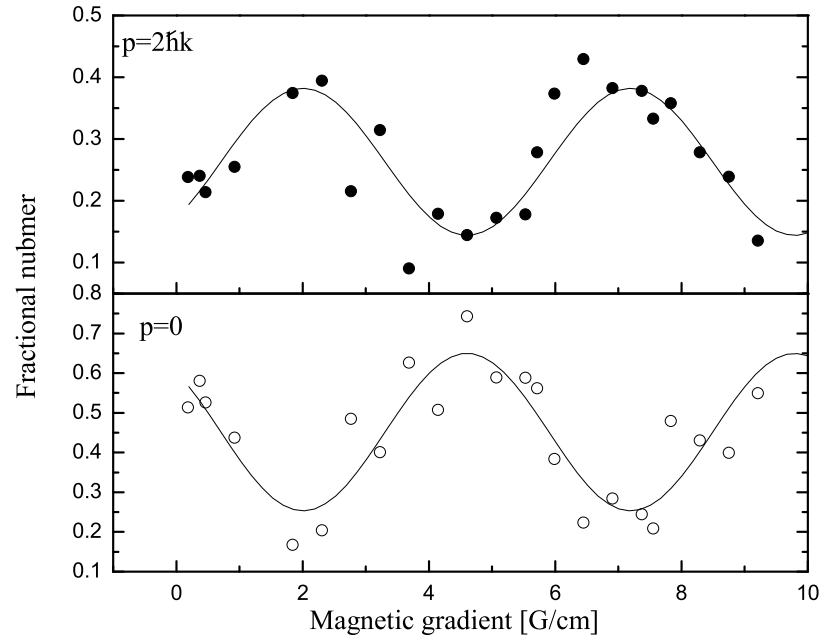


Figure 6.7: Interference after 1 ms propagation with the magnetic pulse turned on at a different time. The average separation of the clouds during the magnetic pulse is $6.94 \mu\text{m}$ when the magnetic pulse is turned on. The expected period is 4.1 G/cm.

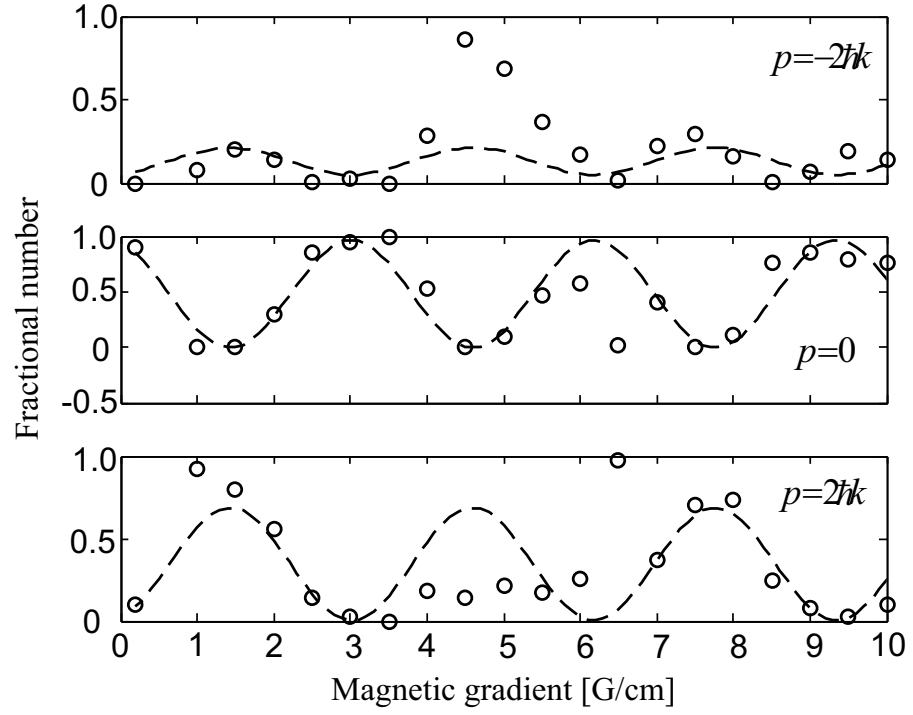


Figure 6.8: Antisymmetric interference patterns of the $|p = 2\hbar k\rangle$ and $|p = -2\hbar k\rangle$ states. The number populations in the $|p = \pm 2\hbar k\rangle$ states have an antisymmetric distribution caused by an initial velocity of the condensate.

two split clouds, and the period of the second interference is twice the normal period.

For an initial velocity of $0.2v_{rec}$, we numerically solve the coupling equations and find that about 10% of the atoms are left in the zero momentum state. Because the residual trap is slightly offset from the magnetic trap that is used to make a condensate, we create a slosh motion if the longitudinal confinement is turned down too rapidly when releasing that atoms from the trap to the waveguide. The theoretical interference curves after the whole sequence are plotted in the Figure 6.9, and the curves agree well with the experimental data.

6.3.2 Collisional heating

Besides the second interference effect, we observe significant collisional heating if the total number of atoms is greater than 2×10^4 . The absorption image in Figure 6.10 shows thermal atoms caused by collision. They are seen in between the peaks of the density distribution after the recombining pulse. The number of collisions per atom can be estimated by a simple formula:

$$N_{collision} \approx \frac{N \times \sigma_{cross}}{\pi \sigma_r^2} \quad (6.8)$$

for a cloud with total number N , radial size σ_r , and s-wave scattering cross section $\sigma_{cross} = 761nm^2$. If we have 20000 atoms with a radial cloud size of $1 \mu m$, the number of collisions per atom is around 0.54. In this case, half of the condensate atoms collide when they pass through each other. The collisions happen when the condensate splits and again when the split clouds recombine. The number of collision decreases when we have fewer atoms and a weaker radial confinement, which gives us a larger radial size. To reduce the collisional heating, we typically use a condensate with a number less than 10^4 and a weak radial trap of 102 Hz.

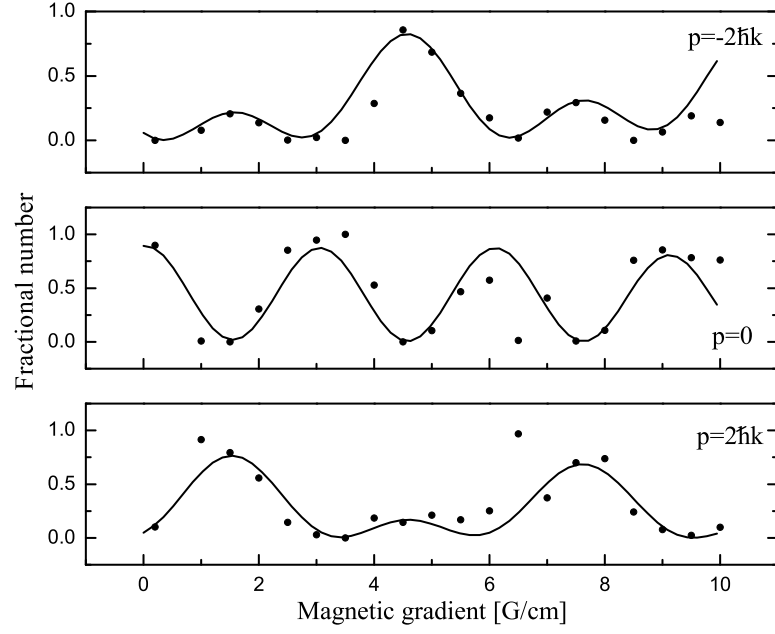


Figure 6.9: Theoretical interference curves of a condensate with a nonzero initial velocity. The solid circles are the experimental data points, as shown in Figure 6.8, and the solid lines are the theoretical curves for an initial velocity of $0.2v_{rec}$.

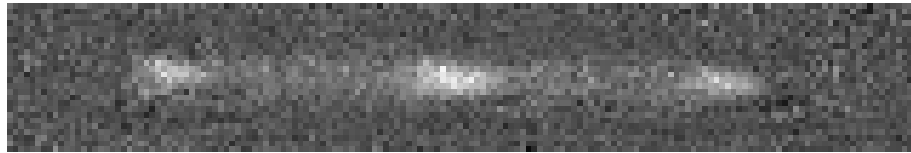


Figure 6.10: Collisional heating. Thermal atoms caused by the collisions are seen in between the peaks of the condensates after the recombining pulse.

6.3.3 Degradation of the contrast ratio

For most of our experimental data, the interference washes out for a propagation time longer than 4 ms. Figure 6.11 shows a plot of the contrast degradation as we increase the propagation time.

The coherence time increases by roughly a factor of two when we do a ‘round-trip’ experiment. Instead of recombining the clouds as they return, we let the clouds pass through each other and apply a second reflecting pulse. The clouds then turn around again and are finally recombined upon their second return to the origin. The interference signals are still observable at 4 ms propagation time with a contrast of around 40%. This suggests that the decay of the contrast ratio is spatially dependent rather than time dependent, possibly due to velocity dispersion caused by the residual trap or the atom-atom interaction.

When the cloud climbs up the trap, different parts of the cloud see different slopes causing different decelerations. A velocity gradient arises across the cloud and so does a phase gradient. When the two clouds interfere upon their return, if the phase gradient is large enough, i.e., greater than π , it creates a rapidly varying density profile because some parts of the cloud interfere destructively while others interfere constructively. After we integrate over the whole cloud to evaluate the number, the interference signals due to the introduced phase shift are averaged out.

On the other hand, an atom-atom interaction also create a phase gradient. We use a condensate with about 10000 atoms. For a 100 Hz radial confinement and a 5 Hz residual longitudinal confinement, the chemical potential is 215 Hz. After the splitting pulse, the atoms experience a higher interaction energy at the place where the the clouds still overlap due to the locally higher density so that their phase evolves faster. A phase gradient is then built up across the cloud. When the two clouds recombine, the atoms that experienced higher interaction energy during the splitting process overlap first,

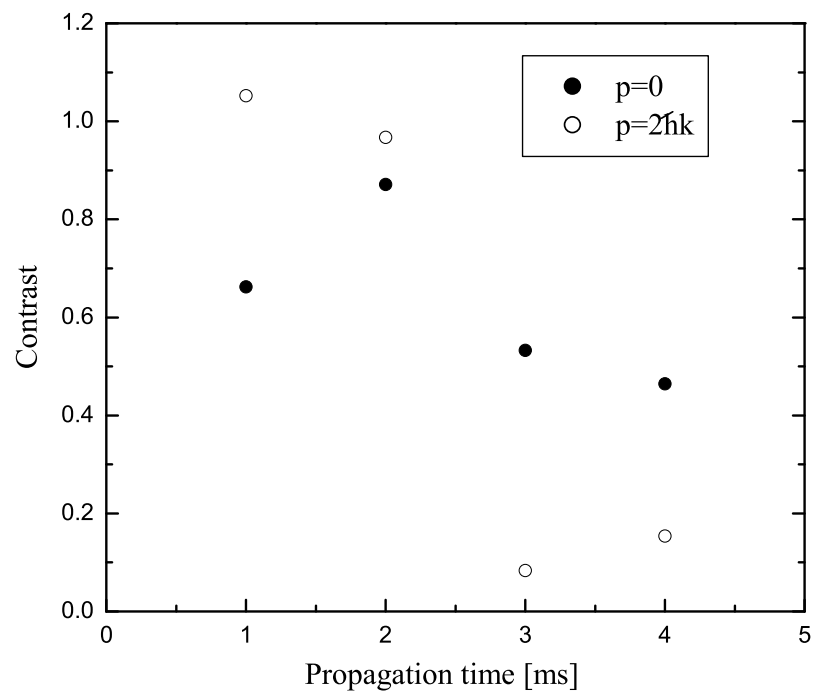


Figure 6.11: Degradation of the contrast ratio. The contrast ratio decays rapidly as we increase the propagation time.

causing more phase gradient to be accumulated. The upper limit of the phase gradient is estimated to be $2\pi \times 215 \text{ Hz} \times 4 \text{ ms} = 1.72\pi$ for a 4 ms propagation time, which is greater than the required π phase gradient.

Some other technical noises such as mechanical vibration and current noise can also cause decoherence. If the standing wave grating shifts half of its period, i.e., 200 nm between the splitting and the recombining pulses, the interference pattern would change from a constructive to a destructive interference. However, we have measured the integrated vibration amplitude to be less than 20 μm by using an accelerometer during the time of our experiment. In addition, our condensate has a lifetime longer than 10 s so the heating caused by the current noise is also negligible on the time scale of the interferometer ($\sim 10 \text{ ms}$).

In summary, the degradation of the contrast could be mainly due to the atom-atom interaction and the residual confinement. A more quantitative study is still required to gain a better understanding and improvement of the contrast ratio for a longer propagation time.

6.3.4 A phase shift introduced by an initial velocity

In the experiment described so far, a magnetic gradient was used to introduce a phase shift and test the interferometer. As an alternative, we can also use the residual 5 Hz trap to create a phase shift. For this, an initial velocity of the condensate is created in the trap with a longitudinal frequency of 5 Hz. The two split clouds have different velocities after applying the splitting pulse. The clouds climb up the potential well to different heights and experience different potentials during the total propagation time, which is equivalent to an effective gradient. As a result, a differential phase shift accumulates as we increase the total propagation time. The atoms are given an initial velocity of 0.445 mm/sec and the differential phase shift is observed by changing the total propagation time in the waveguide. The result of the interference is shown in

Figure 6.12. The best contrast ratio of the interference at 10 ms, with maximum cloud separation of about $120\ \mu\text{m}$, is 20%.

6.4 Prospect for the future work

Our interferometer technique can be used to develop an atom gyroscope. Because the cloud is split in the longitudinal direction, if we can transfer the atoms from a linear waveguide to a ring-shaped waveguide after the condensate is made, the atoms will propagate in the ring waveguide in both clockwise and counterclockwise directions after the splitting pulse. The split clouds will recombine and interfere at their origin after a round trip. The interferometer is thus a gyroscope, which is sensitive to rotation and time-dependent acceleration. With the current linear waveguide design, we can also move the waveguide in time, e.g., up and down, after the condensate splits to make the atoms trajectory form a closed loop to create a atom gyroscope. For 10 ms coherence time, the clouds have a maximum separation of $120\ \mu\text{m}$ so the radius of the gyroscope is $60\ \mu\text{m}$ and the enclosed area is $\pi(60\mu\text{m})^2$. The scaling factor is about $31\ \text{rad}/(\text{rad}/\text{sec})$ and for a π shift, the sensitivity of the rotation is $0.1\ \text{rad}/\text{sec}$. Because the sensitivity of the gyroscope is proportional to the enclosed area, which is proportional to the square of the coherence time, the sensitivity can be increase significantly by increasing the coherence time.

As I mentioned before, the spatially dependent decoherence could be due to the velocity dispersion caused by the residual trapping potential or the atom-atom interaction. To get a longer coherence time, we first need to reduce the inhomogeneity of the potential caused by the imperfection of the wires. Vuletic et al. at MIT has found that the variation of the wires width can be reduced by using a different fabrication process. The variation of the wire with a width of $3.5\ \mu\text{m}$ can be reduced to less than $100\ \text{nm}$ by using the “lift-off” technique and this technique has a tenfold improvement compared to the electroplating process which is currently used in our fabrication. Other

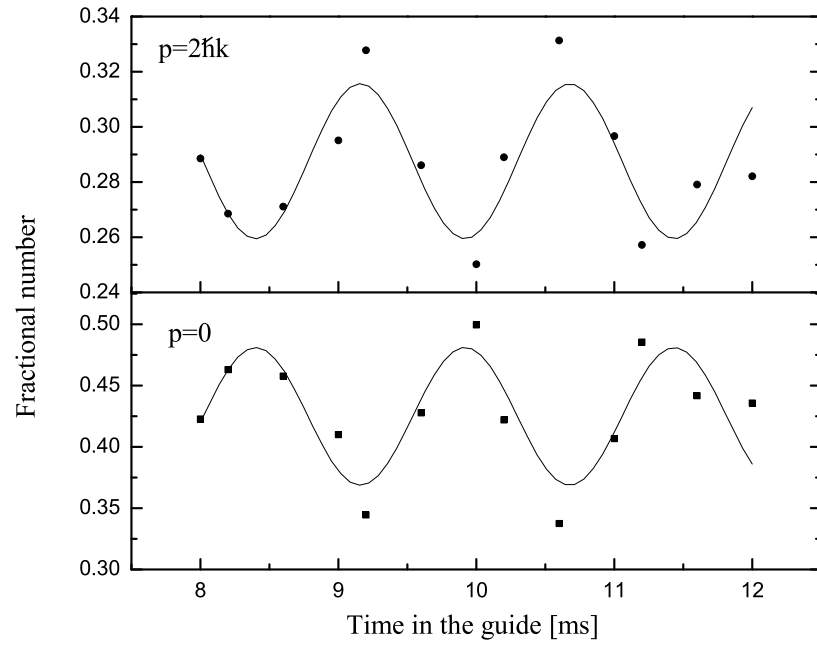


Figure 6.12: Interference after different propagation times in the residual trap. The atoms are given an initial velocity of 0.445 mm/sec to introduce a differential phase shift that develops during the propagation time. The maximum separation of the clouds is 120 μm , at which the clouds are spatially separated.

efforts to improve the quality of the microfabricated wires include using multiple-wire or multilayer designs and making smaller wires to reduce the variation of the wire width in combination with improving the thermal conductivity of the chip to support sufficient current. Second, we need to reduce the atom-atom interaction. One simple way is to reduce the number of atoms in the condensate, however, the signal to noise ratio of the interferometer is also reduced at the same time. Another way is to create a so called “Feshbach resonance”. Magnetically induced Feshbach resonances have been used to alter the scattering properties of the condensates [60, 61, 62, 63, 64]. , where the atom-atom interaction is tuned by a external magnetic field. The Feshbach resonance can also be induced by using optical or rf fields [65, 66, 67]. It has also been proposed that in a waveguide structure, where the atoms can be considered as an one-dimensional system, the atom-atom interaction can be tuned with a strong radial confinement [68]. The presence of the strong radial confinement and the low linear density of the atoms create a “Tonks gas”— a one-dimensional gas of impenetrable bosons, and thus changes the scattering property. For Rb atoms, the one-dimensional scattering length would change sign at a radial confinement of 3 MHz in the Tonks gas regime.

In conclusion, this work reports the first demonstration of an on-chip atom Michelson interferometer employing a Bose-Einstein condensate (BEC). An intra-waveguide optical standing wave serves to split, reflect, and recombine the BEC while the BEC is confined and propagates in a magnetic waveguide. The combined optical beamsplitter and magnetic waveguide employed here have allowed us to study the coherence properties of matter waves confined in a microstructure. The interference is observed up to a round-trip propagation time of 10 ms with the maximum separation of the split wave packets about $120\ \mu\text{m}$. Our technique can be used to make an atom gyroscope with an estimated sensitivity better than 0.1 rad/sec and this sensitivity can be potentially increased by improvement of the coherence time.

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