

Building a Sensitive Atom Interferometer for Navigation

Xiaxi Cheng

Department of Physics
Centre for Cold Matter

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Abstract

The presented thesis describes the setting up of an inertial sensitive atom interferometer for navigation using Rubidium 87 atoms, which is able to perform horizontal acceleration measurement. An ultra-high vacuum interferometer chamber is designed and built. The cold atoms are prepared in a compact 2D MOT setup and loaded into the main MOT. These atoms are then cooled down to few μK and prepared in a magnetically-insensitive ground state using a sequence of optical and microwave pulses, after which atoms are selected in a narrow band of the velocity. A three-pulse Mach-Zehnder ($\pi/2$ - π - $\pi/2$) interferometer using stimulated Raman transitions is used to measure the horizontal acceleration. To improve the sensitivity to acceleration, the fringe contrast of the interferometer requires to be optimised. By investigating the intensity profile of the Raman beams using atoms, a procedure is developed to operate the interferometer in an improved condition, which is able to enhance the fringe contrast by a factor of 3. This procedure improves the shot-to-shot sensitivity of the interferometer below $5 \mu\text{m/s}^2$. In addition, the atom interferometer is compared with a classical state-of-the-art accelerometer.

Declaration

I declare that this thesis is my own work and I have appropriately referenced all sources used in the thesis.

Xiayi Cheng

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Chapter 1

Introduction

1.1 Motivation

In 1978, the first cohort of satellites, collectively known as the Global Navigation Satellite System (GNSS), that enabled users to navigate was launched [1]. Over the last decades, its impressive performance on location tracking has reached an accuracy of few metres [2], which has greatly enhanced transport safety. At the same time, due to economical advantages, today's navigation has become almost inseparable from GNSS. However, the robustness of GNSS can be easily affected by signal jamming and interference. To navigate in an environment, where GNSS is not accessible, alternative guidance system is needed.

An inertial navigation system (INS), as a stand-alone unit, calculates the location of a vehicle from measurements of its acceleration using Newtonian laws without any aid from external position references. Generally speaking, INS measures the inertial forces on three axes using accelerometers and gyroscopes, which provide acceleration and rotation rates, respectively. The early development of INS was mainly driven by autonomous guidance systems for ballistic missile programmes in the 1960s [3], which has led to low cost and much broader civil applications. The performance of such sensors is limited by their bias errors. These contribute to the overall error of dead-reckoned position, which grows quadratically with time. Depending on the purpose of application, these errors might be acceptable for a short mission. However, for long-term applications, a periodic calibration with a reference position (mostly with the aid of GPS) is mandatory. There is great interest now in becoming completely independent of external references, while maintaining a high level of accuracy.

Several technologies have been developed for accelerometers, such as pendulous integrated gyro accelerometers (PIGA) [4], pendulum servo accelerometers [5], vibrating beam accelerometers [6] and accelerometers based on microelectromechanical systems

(MEMS) [7]. Figure 1.1 provides an extensive overview of current applications of accelerometers, which are illustrated as functions of bias and scale factor stability. Due to

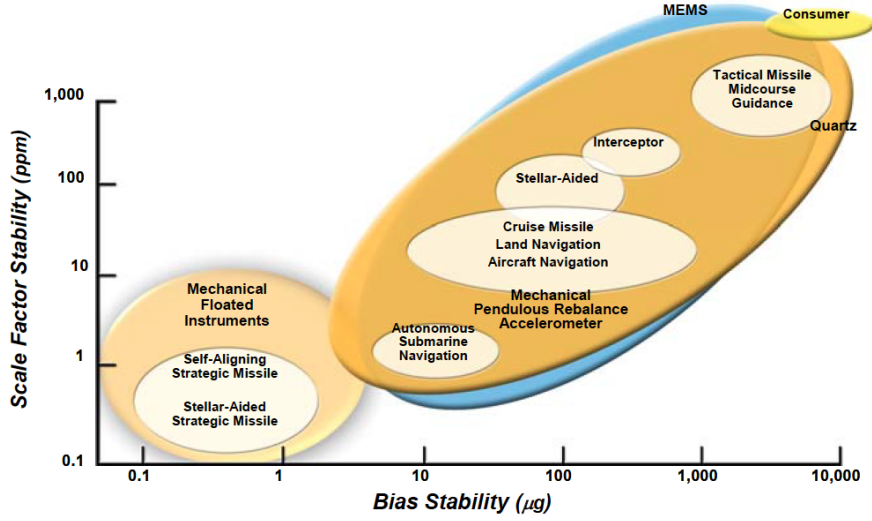


Figure 1.1: Applications of different accelerometer technologies are shown in coloured areas. Bias stability and scale factor stability are used as indicators for their performance ($g = 9.80665 \text{ m/s}^2$ - international standard of gravitational acceleration, which is measured at a latitude of 45° and sea level [8]). Figure taken from Ref. [3].

their low cost and compactness, MEMS and pendulous technologies (Quartz) dominate the majority of the applications.

Despite the high dynamic range and large bandwidth of these accelerometers, they suffer from an unstable bias, which drifts over time. This change can be traced back to variations of the environmental conditions, such as temperature, pressure or mechanical stress in the device. Even in a high-quality device, this causes a positioning uncertainty of few tens of metres¹ after a hour of continuous operation and prevents using INS for long-duration applications. The most promising approach avoiding bias instability is to use cold atom interferometry to measure inertial forces. Cold-atoms, as a near-future INS offer extremely high sensitivity to inertial effects. In addition, as atoms in vacuum are almost isolated from environmental interactions, their atomic properties do not vary during operation and therefore they offer near-ideal stability and reproducibility. The use of sensors based on cold-atom interferometry has great potential for making a stand-alone navigation system without GPS for long-duration missions, which could exceed the current positioning uncertainty by at least an order of magnitude. Figure 1.2 shows the projected level of stability of cold-atom based accelerometers, in comparison with the near-term development of existing technologies.

¹Here, a bias instability of $10 \mu\text{g}$ is assumed.

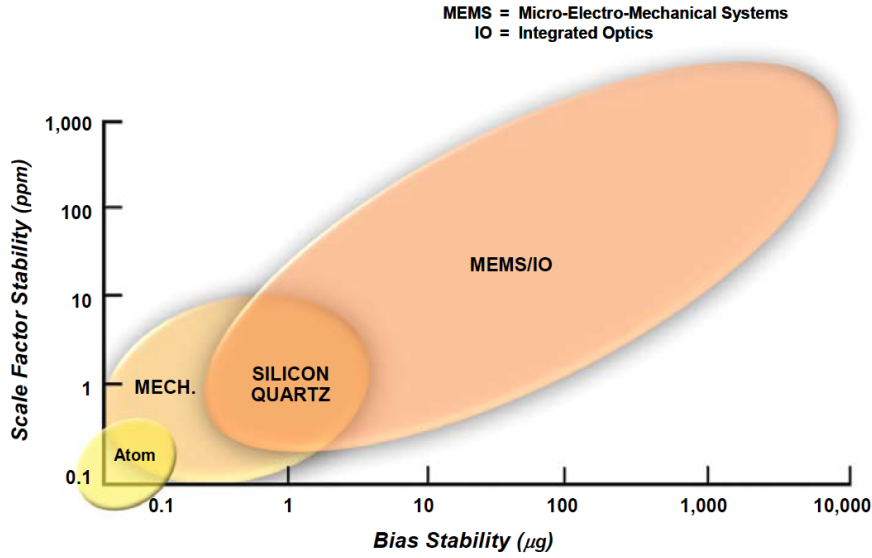


Figure 1.2: Near-term prognosis of performance stability for different accelerometer technologies. MEMS with integrated optics (IO) shows the capacity of improving its performance, to be comparable with silicon and quartz based micro-mechanical systems, whereas cold atom based technology could potentially exceed the stability by an order of magnitude. Figure taken from Ref. [3].

1.2 Inertial sensors using cold atoms

Since the first demonstration of light-pulse atom interferometry by Kasevich and Chu [9], the rapid development of atom interferometry has led it to be the most precise technique for measuring inertial forces, such as acceleration and rotation. Due to its state-of-the-art sensitivity to acceleration of $\Delta g/g \approx 6.7 \times 10^{-12}$ [10], applications of atom interferometry have branched not only into measurement of fundamental constants [11, 12] and testing general relativity [13, 14, 15], but also into development of inertial sensor technology in geophysics [16, 17, 18] and in inertial navigation [19, 20, 21, 22].

The origin of atom interferometry dates back to the development of quantum interference experiments using matter waves. Since Louis de Broglie's postulate [23] on wave-particle duality for any massive object, matter wave interference has been observed with electrons [24], neutrons [25] and atoms [26, 27]. More recently, molecules have also demonstrated their quantum nature with a record mass exceeding 10000 atomic mass unit (amu) [28].

One advantage of using atoms for interferometry lies in the broad range of atomic species (Na, Cs, Mg, He, Ar, Ca, H, K, Li, Ne and Rb) [29], which provide a wide selection of well-known atomic properties, optical scattering cross sections and environmental interactions. Together with the techniques developed in the field of laser cooling and trapping [30, 31], a large ensemble of cold atoms with reduced velocity distribution can be prepared in a small volume in space. This boosts the signal-to-noise ratio during

detection. A narrow velocity distribution also reduces the expansion rate of the cloud. Hence, well-localised atoms can be used for more sensitive interferometry [10, 32], as the output signal of the interferometer scales with the squared of the interrogation time $\propto T^2$.

In light-pulse atom interferometry, Bragg diffraction [33, 34, 35] and stimulated Raman transition [32, 36, 37] are among the most commonly used schemes to coherently manipulate atoms. Generally speaking, Bragg diffraction couples momentum states of the atom without affecting its internal state. The application of higher order Bragg diffractions enables multi-photon recoil and opens up the interferometer area, which, hence, linearly increases the sensitivity of the interferometer [33]. However, due to its high-velocity selectivity, Bragg diffractions are mainly employed in BEC interferometers in fundamental researches. On the other hand, driving stimulated Raman transitions causes excitation between long-lived atomic hyperfine states by 2-photon transition with momentum transfer. This scheme couples the momentum state with the internal state of the atom. Interferometers based on stimulated 2-photon Raman transition also offer the novelty of state-sensitive detection and are able to determine the interferometer phase in a single shot. More importantly, it makes the interferometer fringe contrast insensitive to atom number fluctuations. Experimental imperfections are imprinted either as a phase shift of the fringe pattern or as a reduction of the fringe contrast. Due to the robustness of the interferometry using state-sensitive detection, Raman interferometers have found real-life applications, and they are the subject of this thesis.

Mobile cold-atom inertial sensors

Cold-atom based gravimeters [38, 39, 40, 41], gyroscopes [20, 18, 42] and gravity gradiometers [16, 43, 44, 45] have been realised in laboratory-based setups. As discussed in refs. [46, 47, 48], these atom interferometry devices can be superior in sensitivity and stability to conventional device. However, they are still at a very early stage of development for “in-field” operations, as they typically require well controlled laboratory conditions. Moreover, the highest sensitivities are obtained by allowing the atoms to fall freely for hundreds of milliseconds. Therefore, larger experimental setups are required to accommodate the movement of the atoms.

To move towards a cold-atom based inertial navigation system, the atomic sensors need to be robust, compact and mobile, while maintaining the high level of sensitivity. For this reason, the design of a mobile atom interferometer is distinctively different from lab-based experiments and is accompanied by various engineering challenges. In particular, when operating in noisy environments, the high sensitivity of the atom interferometer can be significantly degraded by the environmental noise.

Several projects have made significant contributions towards this goal in the field of gravimetry. “GAIN” (Gravimetric Atom Interferometer) at Humboldt University of Berlin [48, 49] employs a sophisticated vibration isolation system with a tip-tilt stage to suppress environmental noise and systematic effects due to Earth’s rotation. The “MiniAtom”-project (Miniaturised Coherent Atom Sensors), on the other hand, uses the aid of a classical high-bandwidth accelerometer to overcome the degradation issue introduced by environmental vibration noise [40]. A drastic reduction in physical dimension of the sensing head is achieved by an all-function, hollow pyramidal chamber configuration [17], which unifies the volumes for trapping, interferometry and detection. The success of the “MiniAtom”-project resulted in the formation of μ Quans, which is the second commercial cold-atom interferometry start-up after *AOSense, Inc.*, the first atom interferometry spin-off company from Stanford University. Another transportable atom interferometer is the ICE-project (Coherent atom interferometry for Space). For a short interferometer time, it has proven the capacity to withstand vibrations and temperature fluctuations of a zero-g aeroplane [50].

Towards cold atom based inertial navigation

Despite the commercial success in gravimetry, there is still a substantial void in cold atom technology between the current development and real-life application as an operating stand-alone inertial navigation system, which also involves integration of accelerometers and gyroscopes. A unified interferometry setup for acceleration and rotation measurements [51, 52, 53] has further demonstrated the feasibility of six-axis inertial sensor using cold atoms, which, however, still took place in static laboratory environment for relative low sensitivity. To make the atom interferometer more sensitive to small changes of acceleration, it is desired to have the interferometer operating with a long interrogation time between laser pulses. However, this lowers the bandwidth of the measurement and hence the interferometer doesn’t sense rapid change of acceleration.

Two opposing, yet promising approaches address this challenge, while maintaining high inertial sensitivity of the interferometer in various dynamic environments. One solution is to perform high data-rate interferometry for acceleration measurement with an interrogation time on microsecond scale, which has demonstrated a cycling rate up to few hundreds Hz [52, 54, 55]. However, as the noise of the acceleration of high data-rate interferometry is inversely proportional to the cube of the squared root of the cycling time, $\propto 1/T_{\text{cyc}}^{3/2}$, it is more efficient to perform interferometry with long interrogation time, whereas the acceleration noise scales with $\propto 1/T_{\text{cyc}}^2$. At the same time, improvement on minimising the dead time between measurement cycles would further contribute to interferometer’s performance.

The ultimate goal of our project at Imperial College London is to build a compact, mobile 3-axis accelerometer using atom interferometry, which is suitable for incorporation into a full navigation system. Our cold-atom based accelerometer is required to deliver a measurement noise of $1 \mu\text{m/s}^2/\sqrt{\text{Hz}}$. The project will closely follow the approach of reaching high sensitivity using large T interferometry with a single Raman beam. In combination with a classical, conventional accelerometer, our quantum accelerometer will be able to extend its dynamic range to $\pm 3 \text{ m/s}^2$.

Over the course of my PhD, I have focused on building a transportable, 1-axis accelerometer using atom interferometry with horizontal Raman beams. So far, there is no quantum technology device on the market for measuring horizontal acceleration with the sensitivity offered by vertical gravimeters. Our setup is able to perform cooling, trapping, interferometry and detection on ^{87}Rb atoms in a single chamber. To maintain the simplicity of the setup, we employ a single Raman beam interferometer with a beam waist exceeding 35 mm using 2-inch optics. The Raman optics are fitted into ultra-high vacuum and deliver state-of-the-art surface and wavefront quality, which minimises the systematic error associated with wavefront distortion. To reach the maximal duration for atoms to fall freely across the Raman beams, we employ interferometry with an interrogation time of 50 ms in the horizontal orientation. In order to test its acceleration noise, this quantum accelerometer is compared with a navigation-grade Pendulum Quartz accelerometer. My 1-axis project is to be seen as the precursor to a 3-axis accelerometer. All of the expertise gained will be passed on to the next phase of the project for completing the 3-axis interferometer.

1.3 Thesis outline

The outline of the present thesis follows the natural flow of the experimental progress towards performing a horizontal acceleration measurement aiming for a sensitivity at a scale of $1 \mu\text{m/s}^2/\sqrt{\text{Hz}}$, as illustrated in Fig. 1.3.

Chapter 2 summarises the concept behind the 3-pulse atom interferometry. I start by describing the coherent manipulation of an atom using a stimulated Raman transition. This helps to give a conceptual understanding of the interferometer pulses. I go on to explain how the interferometer measures acceleration. This leads naturally to a discussion of the connection between the measured acceleration and the phase of the lasers driving the Raman transition.

In Chapter 3, I focus on the experimental setup, in which 10^8 atoms are trapped over 100 ms and cooled to 8 μK . After a quick review of the purchased laser systems, I describe the vacuum chamber and other in-vacuum components, together with a home-made light

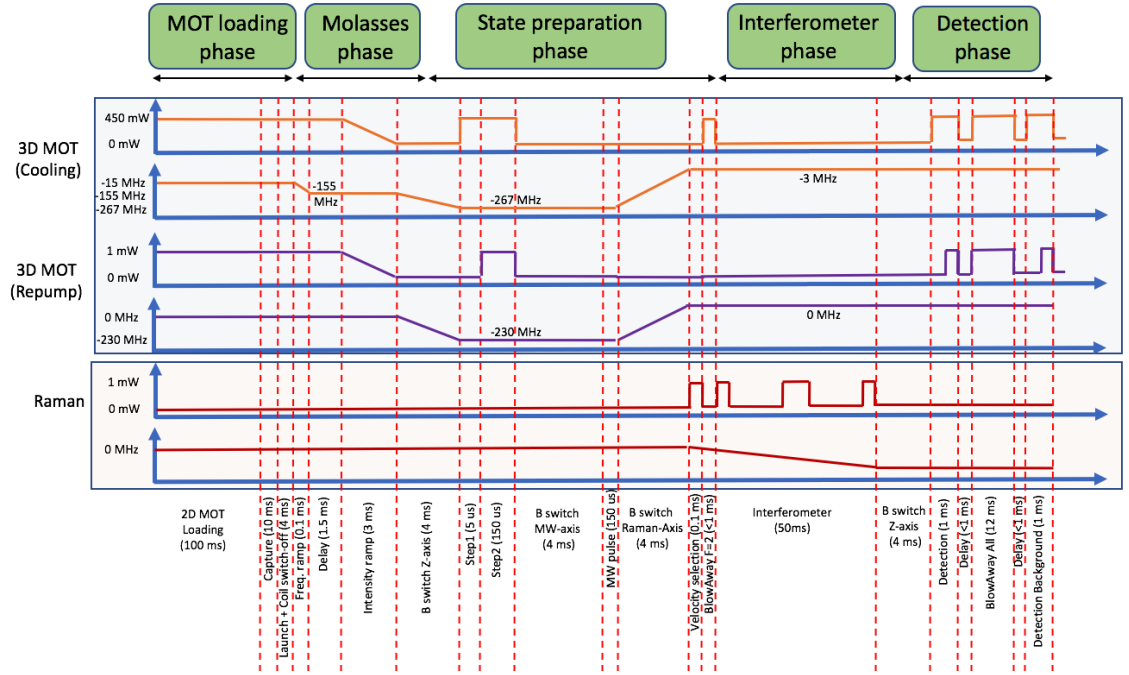


Figure 1.3: Experimental sequence to obtain an acceleration measurement using atom interferometer. This sets the outline of the present thesis.

distribution unit. 2D and 3D MOTs are described individually from their setups to their performance in sections 3.5 and 3.6, respectively. Section 3.7 describes the cooling and launching of the atom cloud using optical molasses, followed by the state preparation scheme outlined in section 3.8.

Chapter 4 gives a more detailed overview of the in-vacuum Raman optics (including the classical pendulum accelerometer) and the lasers that provide the Raman light. My work on understanding and optimising the Raman transition is described in Chapter 5. The characterisation of the interferometer fringes can be found in Chapter 6. Here, I also present the optimisation procedures, which have been carried out to improve the fringe sensitivity by an order of magnitude. Finally, in Chapter 7, the signal of the classical and quantum accelerometers are compared with each other.

Chapter 2

Raman interferometry

Atom interferometry involves coherent manipulation of the de-Broglie wave-packet of an atom. The general design of an interferometer can be divided into 5 steps [29]: state preparation, coherent splitting, free propagation, coherent recombination and detection. The atom interferometer that I build and test in this thesis is directly analogous to the optical Mach-Zehnder interferometer shown in Fig. 2.1. In the optical case, a laser pre-

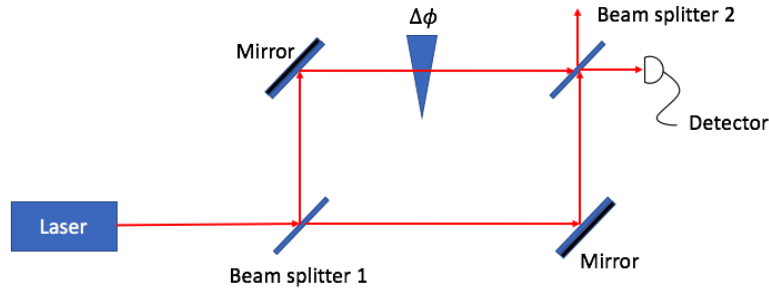


Figure 2.1: A schematic of an optical Mach-Zehnder interferometer. The output of the interferometer depends on the laser phase difference ($\Delta\phi$), which is accumulated between the two paths.

pares a coherent beam (state preparation), which is then split into two separated paths (coherent splitting) by a beam splitter. After some free propagation, the two beams are deflected into a second beam splitter for recombination. We measure the intensity at one of the two outputs. This signal is sensitive to the phase difference ($\Delta\phi$) between the two paths, which is introduced, for example, by an optical wedge extending one optical path length.

My atom interferometer uses a cloud of Rubidium 87 atoms (^{87}Rb), which are prepared in a magneto-optical trap (MOT) and laser-cooled to few μK . Several advantages arise from using laser-cooled atoms. As the atoms are slow, collisions between atoms are infrequent, allowing quantum coherence over long times. As the temperature T is low, the coherence length l_c of the de-Broglie wave packets is long ($l_c = h/\sqrt{m_{\text{Rb}} k_B T}$ [56]),

permitting quantum interference even when the two paths have differing lengths. As there are many atoms with the same interference signal, the signal-to-noise ratio can also be good.

In order to form an interferometer, we require the counterparts of beam splitters and mirrors. We achieve these using stimulated Raman transitions to coherently manipulate the internal quantum states and motional states of the atoms. Section 2.1 summarises the idea of stimulated Raman transition and describes its implementation as beam splitter ($\pi/2$) and mirror (π) pulses. In section 2.2, I describe the working principle of the 3-pulse interferometer constructed from these pulses.

2.1 Stimulated Raman transitions

A stimulated Raman transition is a coherent process, which couples two internal quantum states of the atom. The atom simultaneously absorbs one photon from one laser beam and re-emits it into another laser beam by stimulated emission, which is also accompanied by a momentum transfer.

Figure 2.2 illustrates the D_2 transition ($5^2S_{1/2} \rightarrow 5^2P_{3/2}$) in ^{87}Rb . Over the course of the present thesis, we are interested in driving the hyperfine transition between the magnetically insensitive states ($F = 1, m_F = 0$) and ($F = 2, m_F = 0$), whose energy difference corresponds to the microwave frequency $\omega_0 = 6.834682 \text{ GHz}$ [57]. These are known as the clock states, and I abbreviate them to $|1\rangle$ and $|2\rangle$. Two counter-propagating

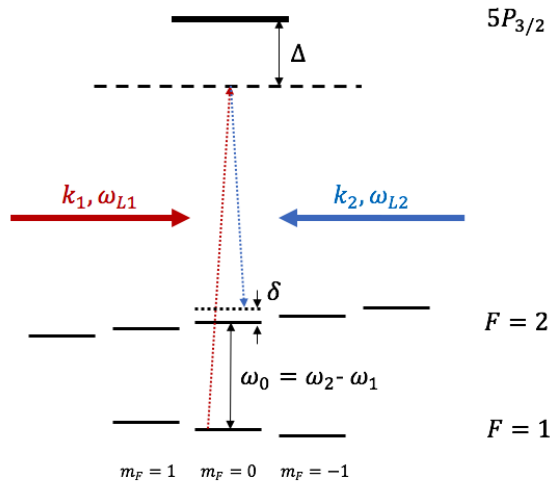


Figure 2.2: Schematic of the 2-photon Raman transition in the D_2 transition of ^{87}Rb .

laser beams with the electric fields described by

$$\begin{aligned}\mathbf{E}_1(x, t) &= \mathbf{E}_{0,1} e^{i(\vec{k}_1 \cdot \vec{x} - \omega_{L1}t + \phi_1)}, \\ \mathbf{E}_2(x, t) &= \mathbf{E}_{0,2} e^{i(\vec{k}_2 \cdot \vec{x} - \omega_{L2}t + \phi_2)},\end{aligned}\tag{2.1}$$

couple the two hyperfine states via a virtual intermediate state, which is noted by the dashed line and is around $\Delta \approx 1$ GHz red-detuned from $5P_{3/2}$ state. Here, $\vec{k}_j$ is the laser wave vector, ω_{Lj} is the optical angular frequency and ϕ_j is the laser offset phase with j specifying the laser beam 1 or 2. I use δ to account for any fine detuning of ω_{L2} from ($F = 2$) to the intermediate state. Since we normally operate with $\delta \ll \omega_0$, it is a good approximation to consider that the laser fields $\mathbf{E}_1$ and $\mathbf{E}_2$ only couple to the transitions $|1\rangle \leftrightarrow |i\rangle$ and $|2\rangle \leftrightarrow |i\rangle$, where $|i\rangle$ denotes the intermediate state. Each coupling is characterised by the Rabi frequency

$$\Omega_j = \frac{\langle j | \mathbf{d} \cdot \mathbf{E}_j | i \rangle}{\hbar},\tag{2.2}$$

where $\mathbf{d}$ is the electric dipole moment operator. The dynamics of the two-photon clock transition is rigorously treated in Ref. [37, 58]. If Δ is sufficiently large that $\Delta \gg |\Omega_1|, |\Omega_2|$ and δ , we can avoid populating the $5P_{3/2}$ state. Then we can eliminate the intermediate state from the equation of motion, and the three-level system is reduced effectively to an analytically solvable two-level system, where the state of the atom can be expressed by

$$|\Psi(t)\rangle = c_1 |1, \vec{p}_0\rangle + c_2 e^{-i(\delta - \phi_{\text{eff}})t} |2, \vec{p}_0 + \hbar \vec{k}_{\text{eff}}\rangle\tag{2.3}$$

with $|1, \vec{p}_0\rangle$ denoting an atom in clock state $|1\rangle$ with momentum $\vec{p}_0$. After the Raman transition to state $|2\rangle$, the atom has momentum $\vec{p}_0 + \hbar \vec{k}_{\text{eff}}$, which has changed because of the photon recoil momentum. The other quantities in Eq. 2.3 are:

$$\begin{aligned}\omega_{\text{eff}} &= \omega_{L1} - \omega_{L2} \\ \vec{k}_{\text{eff}} &= \vec{k}_1 - \vec{k}_2 \\ \phi_{\text{eff}} &= \phi_1 - \phi_2.\end{aligned}\tag{2.4}$$

After solving the Schrödinger equation, the evolution of the coefficients c_1 and c_2 can be described by

$$\begin{pmatrix} c_1(t + t_0) \\ c_2(t + t_0) \end{pmatrix} = U(t, t_0) \cdot \begin{pmatrix} c_1(t_0) \\ c_2(t_0) \end{pmatrix},\tag{2.5}$$

where the propagator $U(t, t_0)$ is

$$U(t, t_0) = e^{\frac{i}{2}\delta t} \begin{pmatrix} \cos\left(\frac{Wt}{2}\right) + i\frac{\delta^{\text{AC}} - \delta}{W} \sin\left(\frac{Wt}{2}\right) & -i\frac{\Omega_{\text{eff}}}{W} \sin\left(\frac{Wt}{2}\right) e^{i(\delta(t+t_0) - \phi_{\text{eff}})} \\ -i\frac{\Omega_{\text{eff}}}{W} \sin\left(\frac{Wt}{2}\right) e^{-i(\delta(t+t_0) - \phi_{\text{eff}})} & e^{-i\delta t} \left(\cos\left(\frac{Wt}{2}\right) - i\frac{\delta^{\text{AC}} - \delta}{W} \sin\left(\frac{Wt}{2}\right) \right) \end{pmatrix} \quad (2.6)$$

and

$$\delta^{\text{AC}} = \frac{|\Omega_2|^2}{4\Delta} - \frac{|\Omega_1|^2}{4\Delta} \quad (2.7)$$

describes the differential ac Stark shifts between $|1\rangle$ and $|2\rangle$ (recall $\delta \ll \Delta$ and we are neglecting the coupling of $\mathbf{E}_1$ to $|2\rangle$ and $\mathbf{E}_2$ to $|1\rangle$). If we start off with all population in $|1\rangle$, $c_1(t_0 = 0) = 1$ and $c_2(t_0 = 0) = 0$, we can determine the probability of finding the atom in $|2\rangle$ as a function of the Raman pulse duration t using Eq. 2.5 and this gives the following expression

$$P_2(p_0, t) = |c_2|^2 = \frac{|\Omega_{\text{eff}}|^2}{W^2} \sin^2(Wt/2), \quad (2.8)$$

where W is defined by

$$W = \sqrt{|\Omega_{\text{eff}}|^2 + (\delta^{\text{AC}} - \delta)^2} \quad (2.9)$$

with

$$\Omega_{\text{eff}} = \frac{\Omega_1 \Omega_2^*}{2\Delta} \quad (2.10)$$

as the effective Rabi frequency. According to Eq. 2.8, the efficiency of driving the atom into $F = 2$ depends on the differential ac Stark shift (or light shift) and the Raman detuning δ . For some values of the detuning Δ in Eq. 2.7, the frequency shift δ^{AC} can be made to vanish by choosing a golden intensity ratio (“the magic ratio”, see section 4.1) for the two Raman laser beams.

The Raman detuning δ , on the other hand, is given by conservation of energy and momentum as

$$\begin{aligned} \delta &= \omega_{\text{eff}} - \left(\omega_0 + \frac{\vec{p}_0 \cdot \vec{k}_{\text{eff}}}{m} + \frac{\hbar k_{\text{eff}}^2}{2m} \right) \\ &= \omega_{\text{eff}} - (\omega_0 + \omega_D + \omega_r), \end{aligned} \quad (2.11)$$

where m is the mass of a ^{87}Rb atom. Equation 2.11 shows that, to be on Raman resonance, the beat note of the Raman lasers ω_{eff} needs to compensate for the Doppler shift ω_D and the recoil shift ω_r , as well as the atomic hyperfine transition frequency ω_0 . Stimulated Raman transitions driven by co-propagating laser beams ($\vec{k}_1 \uparrow \vec{k}_2$) are insensitive to the velocity of the atom, as $\vec{k}_{\text{eff}} = \vec{k}_1 - \vec{k}_2 \approx 0$. These transitions are referred as “Doppler-free” transitions. On the other hand, for beams in the counter-propagating con-

figuration ($\vec{k}_1 \uparrow \downarrow \vec{k}_2$) with $\vec{k}_{\text{eff}} = \vec{k}_1 + \vec{k}_2 \approx 2\vec{k}_1$, the Doppler shift can not be ignored. In section 2.2, we will see that this Doppler shift makes the atom interferometer sensitive to acceleration.

It is worth noting that the Rabi frequency in Eq. 2.2, the effective Rabi frequency Ω_{eff} and the light shift in Eq. 2.7 are as though as the atoms has a single excited state. In our case of ^{87}Rb , a summation over all excited states is required to obtain a more accurate estimate of the Rabi frequency, as analysed elsewhere in more detail [59].

2.1.1 π and $\frac{\pi}{2}$ - pulses

Beam splitters and mirrors for atoms can be realised by driving the stimulated Raman transitions. A beam splitter pulse transforms an atom in a given initial state, $|1, \vec{p}_0\rangle$, into an equal superposition with $|2, \vec{p}_0 + \hbar\vec{k}_{\text{eff}}\rangle$, whereas a mirror pulse reverses the population probability of the atom in either one of these states. Let us demonstrate this by using the propagator in Eq. 2.6. To simplify the notation, we ignore the light shift and assume Raman resonance by setting $\delta = 0$. Equation 2.5 becomes

$$\begin{pmatrix} c_1(t+t_0) \\ c_2(t+t_0) \end{pmatrix} = \begin{pmatrix} \cos\left(\frac{\Omega_{\text{eff}} t}{2}\right) & -i\sin\left(\frac{\Omega_{\text{eff}} t}{2}\right) e^{-i\phi_{\text{eff}}} \\ -i\sin\left(\frac{\Omega_{\text{eff}} t}{2}\right) e^{i\phi_{\text{eff}}} & \cos\left(\frac{\Omega_{\text{eff}} t}{2}\right) \end{pmatrix} \begin{pmatrix} c_1(t_0) \\ c_2(t_0) \end{pmatrix}, \quad (2.12)$$

where ϕ_{eff} is the offset phase difference between the two lasers, given in Eq. 2.4, and Ω_{eff} is the effective 2-photon Rabi frequency given in Eq. 2.10. For simplicity, let us also set $\phi_{\text{eff}} = 0$. For a pulse area $\Omega_{\text{eff}} t = \pi/2$, the input state will be transformed according to

$$\begin{aligned} c_1 |1\rangle &\rightarrow \frac{c_1}{\sqrt{2}} (|1\rangle - i|2\rangle), \\ c_2 |2\rangle &\rightarrow \frac{c_2}{\sqrt{2}} (-i|1\rangle + |2\rangle), \end{aligned} \quad (2.13)$$

where c_1 and c_2 are the probability amplitudes of the respective input states. This pulse is referred to as a $\pi/2$ -pulse. In contrast, when the pulse area $\Omega_{\text{eff}} t = \pi$ (a π -pulse), the probability amplitudes of the input state swap:

$$c_1 |1\rangle + c_2 |2\rangle \rightarrow -i(c_1 |2\rangle + c_2 |1\rangle). \quad (2.14)$$

These pulses allow coupling between different momentum states together with corresponding internal states. Later on, in Chapter 5, I describe and present the effect of Raman interactions on an assembly of cold atom in theory and experiment.

2.1.2 Polarisation requirement

To drive the Raman transition between the clock states, $|F = 1, m_F = 0\rangle = |1\rangle$ and $|F = 2, m_F = 0\rangle = |2\rangle$, one needs suitable polarisations of the two Raman fields. We achieve this by using σ^+/σ^+ circularly polarised laser beams. In the rotating wave approximation, the interaction energy between each σ^+ -polarised Raman laser fields, denoted $\mathbf{E}^{\sigma^+}$, and the atom can be described by

$$\begin{aligned} \mathbf{d} \cdot \mathbf{E}_1^{\sigma^+} &\propto d_1^+ \varepsilon_1^+ e^{-i\omega_{L1} t} + d_1^- \varepsilon_1^- e^{i\omega_{L1} t}, \\ \mathbf{d} \cdot \mathbf{E}_2^{\sigma^+} &\propto d_2^+ \varepsilon_2^+ e^{-i\omega_{L2} t} + d_2^- \varepsilon_2^- e^{i\omega_{L2} t}, \end{aligned} \quad (2.15)$$

where the indices 1 and 2 refer to the corresponding laser fields. I use $\varepsilon_j^\pm$ to describe the amplitude of the electric field $\mathbf{E}_j^{\sigma^\pm}$, which couples to the spherical component of the electric dipole operator, denoted $d_j^\pm$, which induces $\Delta m_F = \pm 1$. The terms with negative exponent in Eq. 2.15 describe the excitation in the dipole transition, and de-excitation by the terms with the positive exponent. The 2-photon transitions between the clock states are, hence, driven by the terms with $d_1^+ d_2^-$ or $d_2^+ d_1^-$ in the 2-photon interaction energy, $\mathbf{d}\mathbf{E}_1^{\sigma^+} \cdot \mathbf{d}\mathbf{E}_2^{\sigma^+}$. This condition is also met when using σ^-/σ^- circularly polarised laser fields.

2.1.3 Beam configuration

When the atom makes a Raman transition, its wavefunction picks up the phase difference, ϕ_{eff} , between the two Raman beams, as shown in Eq. 2.12. Therefore, it is important to have good control of this phase difference, which is achieved by a very high-quality phase lock loop in the laser system (see section 4.1). The two laser beams are delivered to the apparatus in a single optical fibre so that the phase noise arising from fluctuations of the optical path are largely canceled in the difference. The two frequency components are perpendicularly polarised in the fibre, as indicated by H and V in Fig. 2.3. A quarter wave plate (QWP) converts H and V polarised light into beams with handedness of L (left circular) and R (right circular), which drive σ^+ and σ^- transitions, respectively in an atom subjected to a magnetic quantisation field B . These beams are then retro-reflected back through a second QWP by a mirror, which reverses the polarisations. Hence, the retro-reflected k_1 beam drives σ^- transitions and k_2 drives σ^+ transitions. Chapter 4 describes the experimental setup in more detail.

This configuration ensures that the Raman transition is driven by counter-propagating beams, whose polarisation is either σ^+/σ^+ or σ^-/σ^- . These impart a recoil momentum of either $+\hbar k_{\text{eff}}$ or $-\hbar k_{\text{eff}}$, respectively, as illustrated in green and yellow backgrounds

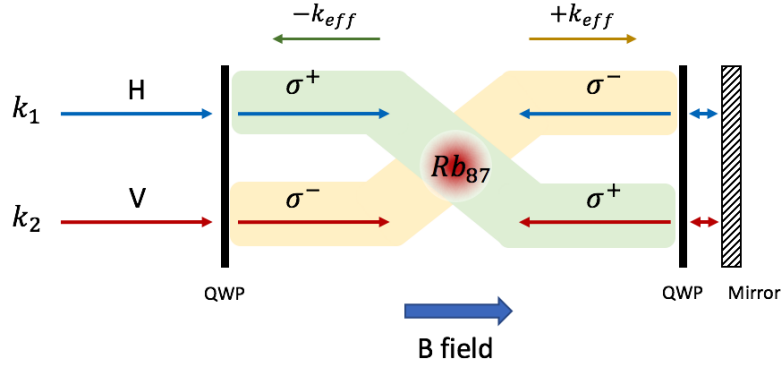


Figure 2.3: Raman beams in the retro-reflecting configuration. Two pairs of Raman beams can be created to drive Raman transition with opposite momentum transfer.

in Fig. 2.3. In the ideal case, when the left QWP is set correctly, this configuration has the advantage that the “Doppler-free” transitions driven by co-propagating beams are suppressed (see Chapter 5 for further discussion of this).

2.2 3-pulse interferometry

Using $\pi/2$ - and π -pulses described in sec. 2.1.1, we can construct a $\frac{\pi}{2} - \pi - \frac{\pi}{2}$ interferometer - a Mach-Zehnder type interferometer. A more detailed description can be found in Ref. [60]. Figure 2.4 illustrates the idea of 3-pulse interferometer in one dimension. An

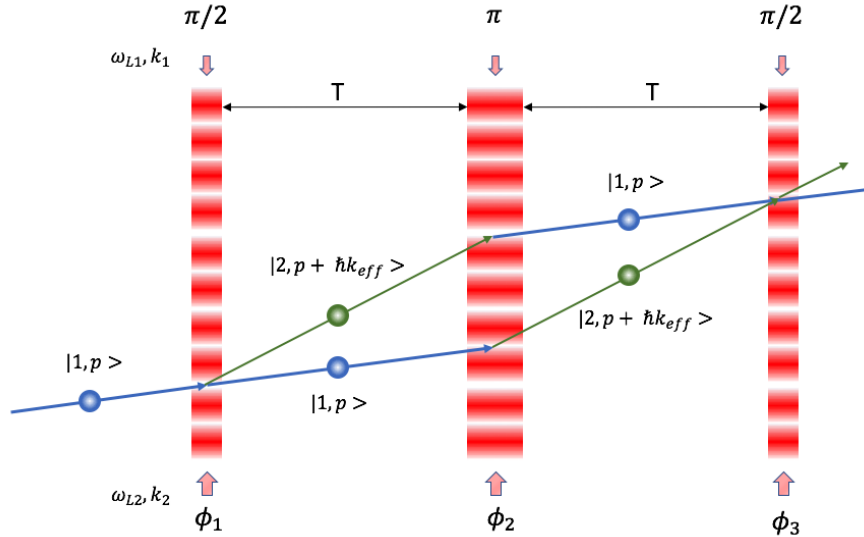


Figure 2.4: Schematic of 3-pulse interferometer: $\pi/2$ - and π -pulses create a Mach-Zehnder type interferometer. The Rabi pulses are based on stimulated Raman transitions, where each Rabi pulse has a well defined phase ϕ_i ($i = 1, 2, 3$). The probability of detecting the atom in $|1\rangle$ and $|2\rangle$ is given by $\cos^2(\Delta\Phi/2)$ and $\sin^2(\Delta\Phi/2)$ respectively, where $\Delta\Phi$ is the phase difference along the two separated paths.

atom with an initial momentum of p is prepared in ($F = 1, m_F = 0$), which is labeled by

$|1, p\rangle$. At $t = 0$, the first $\pi/2$ -pulse splits the atom into a superposition between $|1, p\rangle$ and $|2, p + \hbar k_{\text{eff}}\rangle$. As the atom in $|2\rangle$ has gained an additional momentum $\hbar k_{\text{eff}}$ during the stimulated Raman transition, both wave packets separate in space. After free-propagation of a time T , a second pulse, a π -pulse, is applied, which swaps their states and momenta. After another T of free propagation, the third pulse, a $\pi/2$ -pulse, recombines the wave packets. Using the propagator notation in Eq. 2.5, the state of the atom after the 3-pulse interferometry is given by

$$|\Psi\rangle = U_{\frac{\pi}{2}} P(T) U_{\pi} P(T) U_{\frac{\pi}{2}} |\Psi_0\rangle, \quad (2.16)$$

where $|\Psi_0\rangle$ is the initial state of the atom, $U_{\frac{\pi}{2}}$ and U_{π} are the propagators for a $\pi/2$ and π -pulse respectively. The propagator for the free propagation between the pulses is given by

$$P(T) = \begin{pmatrix} e^{-i\frac{\omega_0}{2}T} & 0 \\ 0 & e^{i\frac{\omega_0}{2}T} \end{pmatrix}. \quad (2.17)$$

The probabilities of detecting the atom in $|1\rangle$ and $|2\rangle$ are determined by the projections $|\langle 1 | \Psi \rangle|^2$ and $|\langle 2 | \Psi \rangle|^2$ respectively. For experimental reasons discussed in section 3.8.1, I define the interferometer signal as the probability P_2

$$P_2 = |\langle 2 | \Psi \rangle|^2 = \sin^2\left(\frac{\Delta\Phi}{2}\right) = \frac{1}{2}(1 - \cos(\Delta\Phi)). \quad (2.18)$$

Here, $\Delta\Phi$ is the phase difference accumulated between the two separated paths, which can be divided in three contributions:

$$\Delta\Phi = \Delta\phi_{\text{prop}} + \Delta\phi_{\text{int}} + \Delta\phi_{\text{light}}. \quad (2.19)$$

The first two of the three are negligible. The propagation phase $\Delta\phi_{\text{prop}}$ is the phase difference that the atom picks up during the free propagation between the “up” and “down” paths and this can be characterised by classical action S_{cl} [60]:

$$\Delta\phi_{\text{prop}} = \frac{S_{\text{cl}}^{\text{up}} - S_{\text{cl}}^{\text{down}}}{\hbar} = \frac{1}{\hbar} \left(\int_{\text{up}} L[x(t), \dot{x}(t)] dt - \int_{\text{down}} L[x(t), \dot{x}(t)] dt \right) \quad (2.20)$$

with $L[x(t), \dot{x}(t)]$ as the Lagrangian for the atom, where x is the position of the atom and $\dot{x}$ the velocity. For a uniform acceleration, which the atom sees, the propagation phase vanishes [60], $\Delta\phi_{\text{prop}} = 0$.

During the free propagation between the interferometer pulses, the phase of the inter-

nal state $|F\rangle$ evolves with

$$\phi_{\text{int}}^F(t_1, t_2) = \int_{t_1}^{t_2} \omega_F dt, \quad (2.21)$$

between t_1 and t_2 . The frequency ω_F specifies the energy of the internal state $|F\rangle$, which varies as a result of the second order Zeeman shift and the ac Stark shift. By stabilising the external magnetic field and Raman laser frequencies and intensities, the interferometer phase contribution $\Delta\phi_{\text{int}}$ due to internal state evolution can be kept as a constant.

We are only left with the phase contribution from the interaction between the laser beams and the atom, $\Delta\phi_{\text{light}}$, and this is the key point for building an accelerometer. As shown in Eq. 2.12, when the atom undergoes the transition from $|1\rangle \rightarrow |2\rangle$, it picks up the phase factor $e^{i\phi_{\text{eff}}}$ and $e^{-i\phi_{\text{eff}}}$ for atoms transferred from $|2\rangle \rightarrow |1\rangle$. Hence, we can summarise all phases accumulated after all three pulses by [60]

$$\Delta\Phi = \Delta\phi_{\text{light}} = \left(\phi_{\text{eff}}^{(1)} - \phi_{\text{eff,up}}^{(2)} \right) - \left(\phi_{\text{eff,down}}^{(2)} - \phi_{\text{eff}}^{(3)} \right), \quad (2.22)$$

where $\phi_{\text{eff}}^{(i)}$ is the phase picked up at i -th laser pulse with $i = 1, 2, 3$. The subscripts “up” and “down” refer to the upper and lower paths in the interferometer.

2.2.1 Acceleration sensitivity

When the atom is subjected to a uniform acceleration $\vec{a}$, the paths of the separate wave packets become asymmetric, as shown in Fig 2.5. For simplicity, let us only restrict the

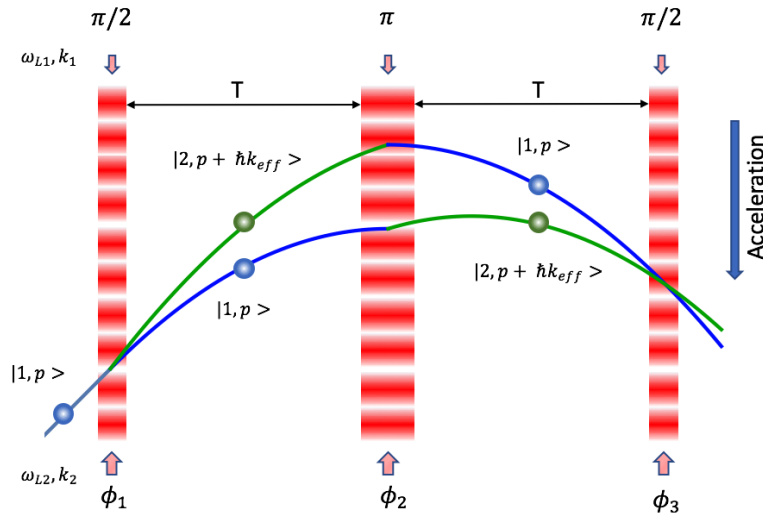


Figure 2.5: $\pi/2 - \pi - \pi/2$ interferometer under acceleration. The linear trajectories of the wave packets are distorted due to acceleration.

acceleration to be along the laser beam axis. In order to ease the notation, we write the

interferometer phase $\Delta\Phi$ in Eq. 2.22 as

$$\Delta\Phi = \Delta\phi_{\text{light}} = (\phi_1 - \phi_2^{\text{up}}) - (\phi_2^{\text{down}} - \phi_3), \quad (2.23)$$

where ϕ_i (for $i = 1, 2, 3$) are the laser phase differences (also referred to as the Raman beat note phase) seen by the atom at all three interferometer pulses. From the frame of reference of the laser, the individual imprinted phase ϕ_i can be expressed by

$$\phi_i = k_{\text{eff}} x(t_i) + \omega_{\text{eff}} t_i + \phi_i^0, \quad (2.24)$$

where ϕ_i^0 is offset phase of the laser, which can be controlled very well in the laboratory. The positions of the atom $x(t_i)$ at the i -th pulse, given at the time t_i , are summarised by

$$\begin{aligned} x(t_1) &= x_0, \\ x_{\text{up}}(t_2) &= x_0 + (v_0 + v_{\text{rec}})T - a\frac{T^2}{2}, \\ x_{\text{down}}(t_2) &= x_0 + v_0 T - a\frac{T^2}{2}, \\ x(t_3) &= x_0 + (2v_0 + v_{\text{rec}})T - 2aT^2 \end{aligned} \quad (2.25)$$

with $v_{\text{rec}} = \hbar k_{\text{eff}}/m$ as the recoil velocity. Substituting Eq. 2.25 and Eq. 2.24 into Eq. 2.23, the final interferometer phase $\Delta\Phi$ can be written as a function of acceleration component a along the Raman beam axis k_{eff} [36]

$$\Delta\Phi = k_{\text{eff}} a T^2 + (\phi_1^0 - 2\phi_2^0 + \phi_3^0). \quad (2.26)$$

From Eq. 2.26, we can conclude that the interferometer phase is determined by the acceleration a , the time T between the interferometer pulses and the effective wave vector k_{eff} . Hence, it is also clear to see that the interferometer driven by co-propagating beams is a factor of $k_{\text{eff}}^{\text{counter}}/k_{\text{eff}}^{\text{co}} \approx 10^5$ less sensitive to acceleration than with the counter-propagating beams.

Apart from the interferometer's sensitivity to acceleration, it is also sensitive to rotation. As the interferometer arms form an enclosed area, any rotation with a component along the normal vector of the enclosed area will introduce an additional phase shift between the two interferometer paths, also referred to as the Sagnac phase. This has also been extensively studied in the field of atom interferometry [18, 42, 60]. Cold atom based gyroscopes will not be covered within the scope of the current thesis.

2.2.2 Fringe sensitivity

Equation 2.18 describes the ideal output signal of the interferometer with a single atom. However, performing Raman transitions on an ensemble of laser-cooled atoms introduces inhomogeneities, which can be imprinted as a reduction of the observed interference fringe contrast C (see Chapter 6). Hence, Eq. 2.18 can be rewritten as

$$P_2 = \frac{1}{2} (1 - C \cos(\Delta\Phi)) \quad (2.27)$$

to describe the measured interferometer signal. The noise of the interferometer can be expressed in terms of inferred interferometer phase noise or acceleration noise, and is determined by the slope of the fringe. The interferometer phase noise $\sigma_{\Delta\Phi}$ is written as

$$\sigma_{\Delta\Phi} = \sigma_{P_2} / \frac{dP_2}{d(\Delta\Phi)} = \frac{2 \sigma_{P_2}}{C \sin(\Delta\Phi)}, \quad (2.28)$$

which is inferred from σ_{P_2} , the noise of P_2 , and the slope of the fringe, $\frac{dP_2}{d(\Delta\Phi)}$. This also indicates that the interferometer is insensitive to the interferometer phase fluctuations on the crests and troughs of the fringe and is most sensitive in the middle, where the sine-term approaches 1. For this reason, the interferometer phase noise is defined at the steepest slope of the fringe. From this, the acceleration noise of the interferometer σ_a can be easily derived and is

$$\sigma_a = \frac{2 \sigma_{P_2}}{C k_{\text{eff}} T^2}. \quad (2.29)$$

Operating the interferometer at a larger T will lower the noise in the acceleration measurement. However, at the same time, the effective operation range Δa of such large T interferometer decreases similarly, as it defined by half of the free spectral range of the fringe (fringe period):

$$\Delta a = \frac{\pi}{k_{\text{eff}} T^2}. \quad (2.30)$$

To lift the degeneracy of interferometer's output, one of the long-term goal of the current project is to use a classical accelerometer with a large operation range to determine the order of the fringe, where the atom interferometer can perform well and correct the drifts in the classical device. Before achieving a hybridised inertial system for navigation, a sensitive atom interferometer needs to be built. This is the focus of the present thesis.

Chapter 3

Preparation of the atoms

In order to develop a cold-atom accelerometer, the atoms need to be in ultra-high vacuum and cold enough to interact with the Raman light over some tens of milliseconds. A large collection of atoms is required to obtain measurement with a high signal-to-noise ratio. To satisfy these requirements, laser cooled atoms are the ideal candidates.

Our system is based on Rubidium-87 (^{87}Rb) atoms. In this chapter, I describe the experimental setup that we use to control, to trap and to laser-cool Rb atoms. The assembly of atoms is then prepared in the lower clock state, which is the starting state for the interferometry.

3.1 Device control

In order to build a working device from an empty table, the structure and performance of the hardware and the development of the software are inseparable from each other. A wise choice of experimental hardware can significantly simplify the structure of an experiment, which leads to elegant system control and, additionally, opens up opportunities for further development with simple system upgrade. Moreover, thoughtfully written software guarantees reliable communication between hardware components, which is essential for good performance of the device.

The device control presented in this section shows the current, most updated control structure, which has been iteratively improved over the course of the project.

3.1.1 Hardware

Figure 3.1 illustrates the structure of the hardware components for the interferometer. We require precise control of the magnetic fields and the intensities and frequencies of the lasers and the microwave. The user sends analog and digital signals via *National Instru-*

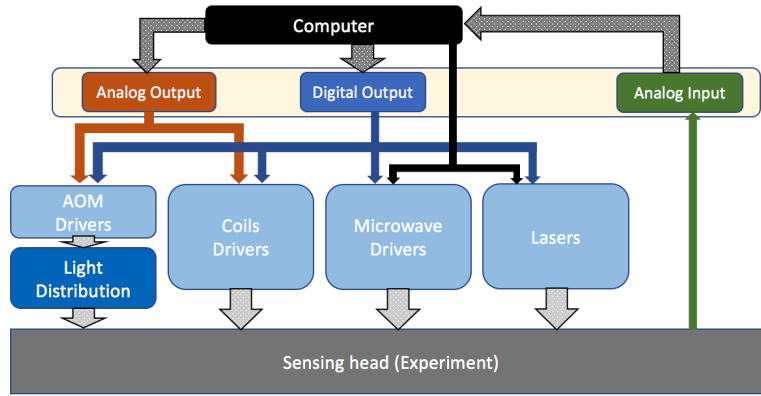


Figure 3.1: Schematic of the hardware structure. Analog signals (orange) control the AOM frequencies and magnetic fields. Serial commands (black) are sent to vary the microwave and laser frequencies. Digital signals (blue) are used to trigger all components, whereas the detected signal (green) is fed back to the computer via a 24-Bit DAQ-card.

ments modules (*PXI 6732* and *PXI 6541*) to individual components. The analog signals, as shown in orange arrows in Fig. 3.1, are assigned to control the frequencies and amplitudes of six rf signals individually, and the currents going into six separated magnetic coils. The rf signals are used to drive the acoustic-optical modulators (AOMs) for more control over the MOT laser beams (see section 3.4), which are then guided into the sensing head. More descriptions on the magnetic coils can be found in section 3.6.1. The digital signals, shown as blue arrows in Fig. 3.1 and originating from a 50 MHz clocked *National Instruments* module (*NI-HSDIO*), are triggers to switch on and off the components precisely on demand. Serial connections shown in black, pre-set the frequencies and amplitudes of the microwave source (see section 3.8.3) and the laser source (see section 3.2). Finally, the detection signal (green arrow) is acquired by a 24-Bit DAQ-card (*PXI 4462*) from *National Instruments*, which is further used for signal processing on the computer.

3.1.2 Software

The control software is called *MOTMaster*, and is based on a collection of low-level C# programmes, which were developed within the Centre for Cold Matter at Imperial College London. In the initial stage of the project, when *MOTMaster* was not operational, we used MIT’s freeware *Cicero Word Generator* [61], which has inspired *MOTMaster*’s user interface (UI) development.

The overall design of *MOTMaster* relies on performing tasks using low-level functions to minimise computational overhead. A user interface allows the user to design and build sequences of tasks. Figure 3.2 shows *MOTMaster*’s basic structure. The *SequenceData* class acquires data from the user interface. These data are used by *SequenceIO*

class to generate patterns for each object defined in the *NavigatorHardware* class, such as the analog and digital hardware channels, the RS232 serial connection (with μ Quans lasers) and the connection to ICEBloc (M-Squared lasers). The patterns represent the temporal changes of an output channel and simplify synchronisation between multiple channels. The execution of the output patterns are achieved then by *Controller* class. The *ExternalControl* class enables remote access of *MOTMaster* by other softwares, which is used to visualise data. An extensive description of *MOTMaster* can be found in Ref. [62].

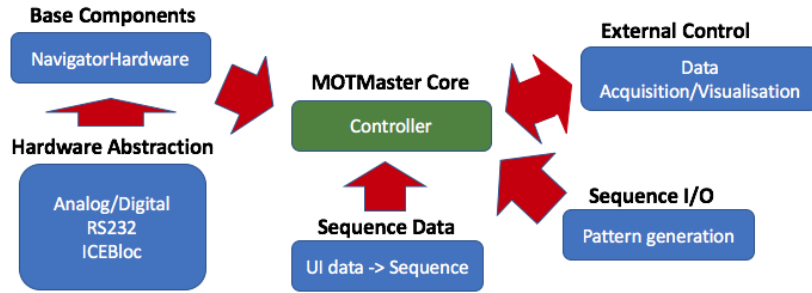


Figure 3.2: MOTMaster’s basic structure. User-specified sequence data are prepared as hardware output patterns. By defining each hardware channel as an object, *Controller* executes the pattern on the corresponding channels. An external software communicates with the MOTMaster’s *Controller* to visualise the data.

3.2 Laser systems

Two independent laser systems are used. One system, referred to as the MOT lasers, addresses the cooling and repump transitions of ^{87}Rb on the D2-line. This laser system is not only responsible for laser cooling and state preparation, but also for state sensitive detection and the blow-away beam, all shown in red in Fig. 3.3. The second, more powerful system is dedicated to driving stimulated Raman transitions between the magnetically insensitive hyperfine ground states in $5^2S_{1/2}$ level and is referred to as the Raman lasers, which are described in more detail in Chapter 4. In Fig. 3.3, the Raman laser frequencies are illustrated by purple lines.

3.2.1 MOT lasers - μ Quans

The MOT laser system was purchased from μ Quans. It gives out two frequencies in the same polarisation, overlapped together in a single fibre output. This is used in combination with a light distribution unit (see section 3.4) to provide multiple outputs and frequencies for different stages of the atom interferometry.

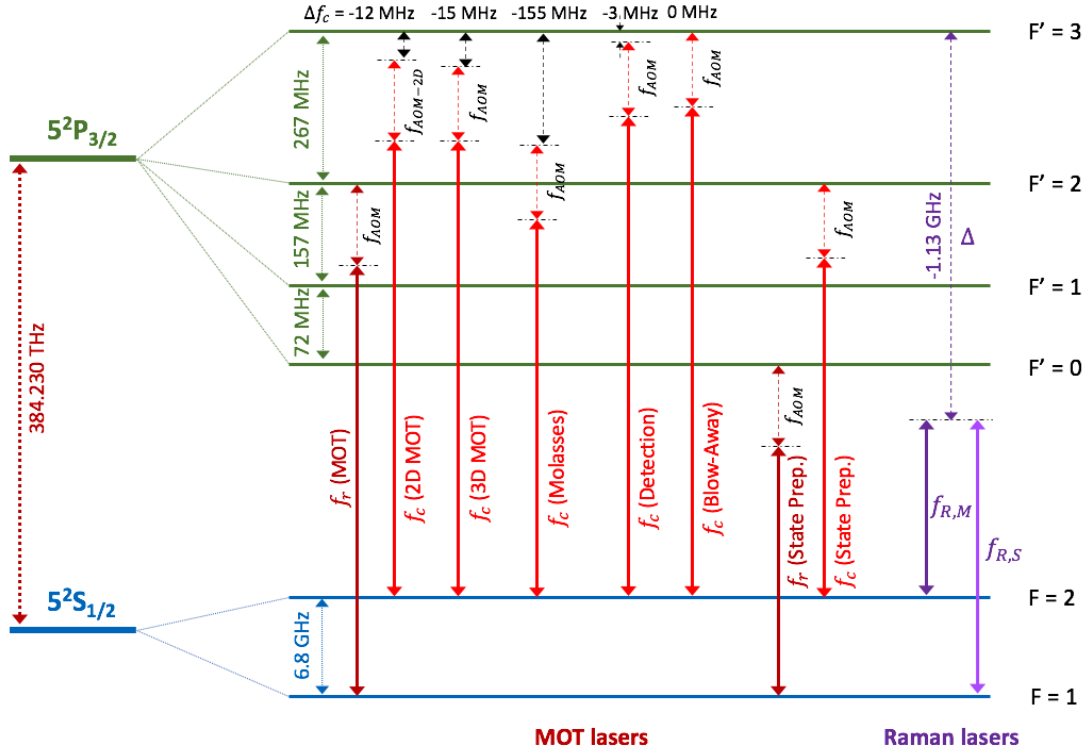


Figure 3.3: Rubidium 87 D2 line with the MOT cooling and repump frequencies, denoted by f_c and f_r respectively, and the Raman master and slave laser frequencies by $f_{R,M}$ and $f_{R,S}$. All cooling and repump frequencies are reached using external fibre AOMs to produce the required detunings at various stages of the experiment.

A reference laser in the MOT laser system is locked on to atomic transitions of ^{85}Rb atoms. This reference frequency is used in an offset-lock to stabilise the output frequency f_c of a slave laser. Before reaching the atoms, the slave laser beam f_c goes through a fibre acoustic-optical modulator (AOM) in the light distribution unit, which increases the slave laser frequency f_c by the carrier frequency of the AOMs, f_{AOM} , as indicated by the dashed red arrows in Fig. 3.3. The output frequency $f_c + f_{AOM}$ is locked closely to the cooling transition for a magneto-optical trap (MOT), $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transition in the D2 line of ^{87}Rb . (Here, $|F\rangle$ and $|F'\rangle$ indicate the hyperfine structure $5^2S_{1/2}$ and $5^2P_{3/2}$ states, respectively.) By precisely controlling the offset frequency, we can vary the absolute detuning Δf_c of the laser frequency, $f_c + f_{AOM}$, from the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ resonance. This allows us to perform (sub-)Doppler cooling [30] (see section 3.7), state preparation (see section 3.8), detection (see section 3.8.1) and blowaway. Figure 3.3 exhibits the absolute detunings in black dashed arrows of the cooling light frequencies of the MOT lasers for various stages of the atom interferometry.

The cooling beam also off-resonantly addresses $|F = 2\rangle \leftrightarrow |F' = 2\rangle$ transition. From ($F' = 2$), atoms can decay into ($F = 1$) and become dark to the cooling light. To put all lost $F = 1$ atoms back into cooling cycle, a sideband with a frequency f_r is

generated from the cooling beam by an electro-optical modulator (EOM). Together with the fibre AOM in the light distribution unit, the overall frequency, $f_r + f_{\text{AOM}}$, is always kept on resonance with $|F = 1\rangle \leftrightarrow |F' = 2\rangle$ transition throughout (sub) Doppler cooling and detection. This is achieved by adjusting the driving modulation frequency of the EOM to counter-balance the frequency change of the cooling frequency. These two beams are finally overlapped together in the same linear polarisation and delivered in a single mode polarisation maintaining fibre (PM fibre). We measured a maximum power of 550 mW with 1 % power fluctuation and the worst polarisation extinction ratio (PER) of -20 dB.

3.3 Vacuum chambers

Figure 3.4 (a) shows a schematic of our vacuum chambers. Two chambers, a 2D^+ MOT source chamber and a 3D MOT interferometer chamber, are connected via a DN16 vacuum valve for modular vacuum diagnostics. The 2D MOT source chamber is a UHV

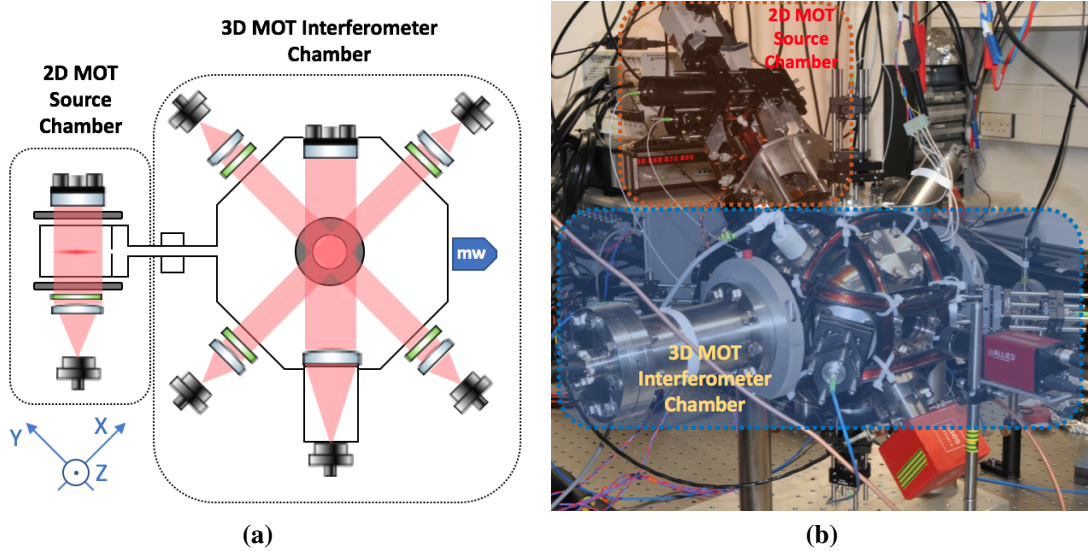


Figure 3.4: Schematic of the experimental setup (a). A 2D MOT is formed in a secondary chamber and fills a six-beam 3D MOT at a higher rate. Laser cooling and interferometry take place in the 3D MOT interferometer chamber. The Raman optics are situated inside the vacuum and aligned with the microwave horn (mw) in the X/Y -plane. A picture of the setup is shown in (b).

glass cell, purchased from *Cold Quanta Inc.* The cell has the dimension of 30 mm x 30 mm x 93 mm. Towards one end, a reflective silica wafer with an aperture of a diameter of 0.7 mm is mounted in the cell allow a high Rb pressure in the 2D^+ MOT cell without imposing an undue gas load to the 3D MOT chamber. Two Rb dispensers are already built inside the cell with electrical connectors. The cell is bonded to a DN16CF flange,

attached to the vacuum valve. The optical assembly for the $2D^+$ MOT is built as a module, fitted closely around the glass cell. The purpose of using a $2D^+$ MOT is to achieve a high MOT loading rate in the 3D MOT interferometer chamber, which, in turn, reduces the dead time in the interferometer sequence, while maintaining a very low pressure in the main chamber. Section 3.5 describes the optical setup of the 2D MOT chamber, its performance and operation in more detail.

Figure 3.4 (b) shows a photograph of the 3D MOT interferometer chamber with all flanges, windows, valves and coils attached. The chamber is a bespoke spherical cube made by *Kimball Physics* using 316L stainless steel with a magnetic permeability very close to 1. This results in an almost completely non-magnetic chamber, helping to minimise undesirable stray magnetic fields, which could distort the cold atom cloud. The inner and outer diameters of the chamber are 16.51 cm and 20.7 cm, respectively. This provides enough surface area to situate 16 x DN40CF and 2 x DN63CF flanges. I use six of the DN40CF flanges to provide optical access for the MOT beams, one for the ion pump, one for the microwave horn, one for imaging of the cold atoms and one for DN40CF to DN16CF adapter flange, which is required to connect with the 2D MOT source chamber. The remaining DN40CF flanges allow access for further diagnostics, should they be required in the future. The two larger DN63CF flanges are used to fit the Raman optics into the vacuum. One of the DN63CF flanges is extended by 15 cm to accommodate the Raman beam collimator, while the other one is occupied by a high-quality retro-reflecting mirror and a quartz-pendulus accelerometer attached to the back of the mirror. More detailed descriptions of the components inside the vacuum are given in section 3.6 and chapter 4. To ensure that this chamber could achieve ultra-high vacuum, we first cleaned each component using a series of solvents (water, acetone and isopropanol). Next, individual components were baked according to their maximal temperature ratings in order to outgas any remaining solvents. All components are carefully chosen for a low outgassing rate ($\leq 10^{-8}$ mbar l/s) and tested in a separate test chamber.

To get to UHV, both chambers, flanges, viewports and ion pump were fixed together and evacuated using a turbo molecular pump. This was then baked as an assembled system for a further 48 hours to remove further contaminants. Next, the 100 l/s ion getter pump, (*SAES getters NEXTorr D 100-5*) was switched on. A valve was used to separate the turbo molecular pump from the rest of the system before it was switched off and removed. After one week of pumping under the sole action of the ion getter pump, our clean chamber achieved a base pressure of 10^{-9} mbar. In order to put vacuum coils and vacuum Raman optics inside, we vented the chamber using Dry Argon N6.0 at the highest grade of purity. Once the vacuum was under atmospheric pressure, an over-pressure was created before opening and inserting the new components. We repeated the pumping

process and reached a base pressure of $6 \cdot 10^{-10}$ mbar after a month.

From this point onwards, I define a cartesian coordinate system with gravity pointing along the Z axis, for convenience. The chamber assembly is oriented so that each counter-propagating pair of MOT beams is aligned with a cartesian axis, as shown in Fig. 3.4 (a). The Raman beams are located on the X/Y -plane (perpendicular to the gravity for horizontal acceleration measurement) and at an angle of 45° to X - and Y -axes.

3.4 Light distribution

The Raman (*M-Squared*) and MOT (*μ Quans*) lasers provide us with one output fibre each, as described in section 3.2. While the Raman light is ready to use, with no further manipulation of its frequency or intensity required, the MOT light is more complicated to generate, requiring 8 separate beams, each of which requires individually precise control over the frequency and power. Six of them are dedicated to the three pairs of counter-propagating MOT beams, labeled with (X_1, X_2) , (Y_1, Y_2) and (Z_+, Z_-) . The remaining two, used to generate the $2D^+$ MOT and to push the $2D^+$ MOT into the main chamber, are referred to as the 2D MOT beam and the push beam respectively.

We distribute the MOT light using optical fibres throughout to ensure both compactness and robustness. Figure 3.5 shows a schematic of the light distribution unit used, which is based on a series of customised fibre beam splitters (purchased from *Evanescent Optics Inc.*) and fibre AOMs (purchased from *Gooch & Housego*). The ratios for

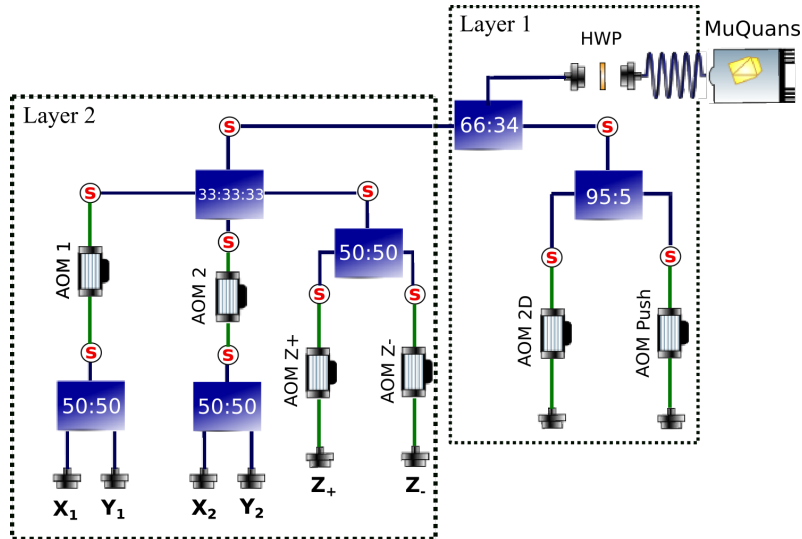


Figure 3.5: Layout of the fibre distribution unit: Layer 1 includes the fibre optics for the 2D MOT, two fibre couplers and a zero order half-wave-plate (HWP). Layer 2 includes all fibre splitters and fibre AOMs for the 3D MOT. All connections between fibres optics are spliced together and labeled by a red "s".

power splitting are chosen so as to optimise the power in each beam, with the intrinsic

loss of each component taken into account.

The distribution unit is divided into two stacked layers. The first layer contains a free-space fibre coupling unit with a half-wave-plate to have better control over the polarisation alignment into the input of the distribution unit. In addition, the fibre optics responsible for generating the $2D^+$ MOT is also hosted on this layer. The fraction of the total power going to the $2D^+$ MOT distribution unit is 34 %. Of this, 95 % is used for cooling the atoms and 5% for the push beam. Using fibre AOMs, we optimise the frequency and power individually.

For the 3D MOT, 66 % of the total power passes through a 33% : 33% : 33%-splitter. We then guide the light from two 33 %-outputs into separate fibre AOMs (AOM 1 and AOM 2), before splitting them again equally into two. These four beams (X_1, X_2, Y_1, Y_2) create the cooling beams along the X and Y -axes and enable launching of the molasses along the Raman axis. Section 3.7.1 describes launching in more detail. The beams labeled with $Z_{\pm}$ are generated by the light originated from the third 33 %-output, which is first split and then sent through two separate fibre AOMs. They are equipped with individual power and frequency control by the AOMs. We use Z -beams not only for the 3D MOT, but also for launching the atoms against the gravity, as imaging beams (see section 3.8.1), blow-away beam, and to perform purely optical state preparation, as described in section 3.8.2.

The first implementation

In my first implementation of the distribution unit, I connected each fibre using mating sleeves. These are threaded metal tubes, which are supposed to press two fibre ends tightly against one another to ensure that their principal axes are aligned with a precision better than 1 degree. I tried using the best commercially available mating sleeves. However, I found that each mating sleeve has, on average a 20 % power loss. This becomes more severe at higher input powers. In addition to the inefficient coupling, each mating sleeve degrades the polarisation purity. The observed polarisation extinction ratio (PER) change for each coupler is about 10 dB, which results in a variation of the splitting ratio up to 40%. The resulting power and polarisation drifts cause a highly unstable MOT with a fluctuating atom number.

Splicing

To overcome the problems associated with the mating sleeves, I used an *Ericson FS* fusion splicer to make the joins. This was a delicate and time consuming procedure. First, each fibre end was cleaned and cleaved. Then, using a microscope the slow axes of the two fibres to be spliced were aligned to one another with a precision of less than

1 degree. In order to achieve higher power transmission and small PER degradation, I manually performed polarisation alignment in-situ by maximising the PER and the power transmission during the splicing process. Next, the tip of each fibre was melted and the fibres were joined together into a single piece. The power loss at each spliced junction was measured to be $\leq 5\%$, which was able to double the laser power, on average to 20 mW, measured at the output of each MOT fibre before going into a MOT collimator (see sec. 3.6.1). In addition, splicing improved the loss of PER at each joint from 10 dB to ≤ 0.5 dB. Among all outputs of the MOT fibres, the worst PER measured is -15 dB, which is limited by the fibre AOMs and lies within the specifications provided by the manufacturer.

3.5 $2D^+$ -MOT

Our interferometer aims to operate within an interferometer time of $2T = 50$ ms. For this to operate at its highest sensitivity, we want the quantum projection noise, as the fundamental noise, to dominate amongst other systematic and statistical noise sources. The uncertainty in acceleration σ_a (see section 7.1.3) per measurement is then inversely proportional to $\sqrt{N}$, where N is the number of atoms detected in the interferometer. On the other hand, a larger atom number requires a longer loading time. When loaded from the background vapour pressure, it can take up to a few seconds to fully load the MOT. This significantly contributes to the dead time T_{dead} in sequential operation of the atom interferometer and decreases the data rate. Here, I define the data rate as the cycling frequency of the interferometer, $f_c = 1/(2T + T_{\text{dead}})$, where the dead time, $T_{\text{dead}} = t_{\text{load}} + t_{\text{prep}} + t_{\text{det}}$, consists of the MOT loading time, the state preparation time and the detection time. As t_{load} has the largest contribution, I approximate $T_{\text{dead}} \approx t_{\text{load}}$. Figure 3.6 exhibits the acceleration uncertainty σ_a of a $2T = 50$ ms interferometer after 1 second of measurement, which is normalised against the fringe period of the interferometer, $\Delta a = \pi/(k_{\text{eff}} T^2)$. This fractional acceleration uncertainty, $\sigma_a/\Delta a$ is plotted as a function of the dead time, T_{dead} . Red, black and green lines indicate the fractional acceleration uncertainty for different MOT loading rates of $\gamma = 1/t_{\text{load}} = 1, 10$ and 100 s^{-1} , respectively. The growth of the atom number in the interferometer, $N_{\text{int}}(t)$, is proportional to the growth of the atom number in the MOT [63], and is given by

$$N_{\text{int}}(t) = r_{\text{sel}} N_{\infty}^{\text{MOT}} (1 - e^{-t\gamma}), \quad (3.1)$$

where r_{sel} is the portion of the MOT atoms selected for the interferometry (see sec. 5.1) and is taken to be 1/10 for the calculation in Fig 3.6. N_{∞}^{MOT} is the equilibrium atom

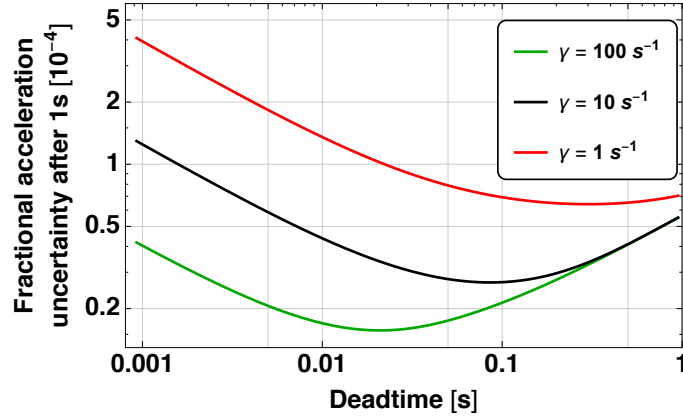


Figure 3.6: Fractional noise $\sigma_a/\Delta a$ of a $2T = 50$ ms interferometer after 1 s of measurement is shown for MOT loading rates of $\gamma = 1, 10$ and 100 s^{-1} . The uncertainty initially decreases as the atom loading time increases, but then it increases with stagnating number of atom due to decreasing data rate. An optimal operation condition occurs for the dead time to be close to the loading time $t_{\text{load}} = 1/\gamma$.

number in the 3D MOT and assumed to be 10^9 to approximate our experimental condition. Two counter-acting trends are visible here. At shorter dead time, while the atom number in the MOT increases, the acceleration uncertainty decreases due to decreasing quantum projection noise. As the atom number stagnates at longer dead time, the data rate decreases, which results in a higher acceleration uncertainty. This is shown in the rising lines on the right hand side of the figure. From this, one can conclude that an ideal trade-off between atom number and data rate can be reached, when the dead time is close to the MOT loading time, $1/\gamma$.

For this reason, we shorten the interferometer's dead time in our setup by using a rapid loading 2-dimensional magneto-optical trap (2D MOT) [64]. In a 2D MOT, atoms are optically cooled on two axes and form a cigar-shaped MOT along the field-free axis of a two dimensional magnetic quadrupole. As the 2D MOT forms in a small glass cell with high Rb partial pressure (separated from the main chamber), the loading is very fast. Another slightly red-detuned beam is aligned along the non-confining axis and retro-reflected back by the silicon wafer with an aperture in the glass cell. The retro-reflecting beam introduces additional cooling, which decreases the longitudinal velocity of the atoms. For this reason, this configuration is also referred as a $2D^+$ -MOT. Moreover, the retro-reflecting beam also acts as a push beam to launch atoms into the 3D MOT chamber through a 0.7 mm-diameter aperture. Based on the existing 2D MOT glass cell, I have designed and built a compact, modular setup to unify the optics and magnets for the $2D^+$ -MOT. This system has demonstrated a 3D MOT loading rate of 10 s^{-1} and it can comfortably load 10^8 atoms in ≤ 100 ms with 65 mW of total optical input power.

3.5.1 Setup

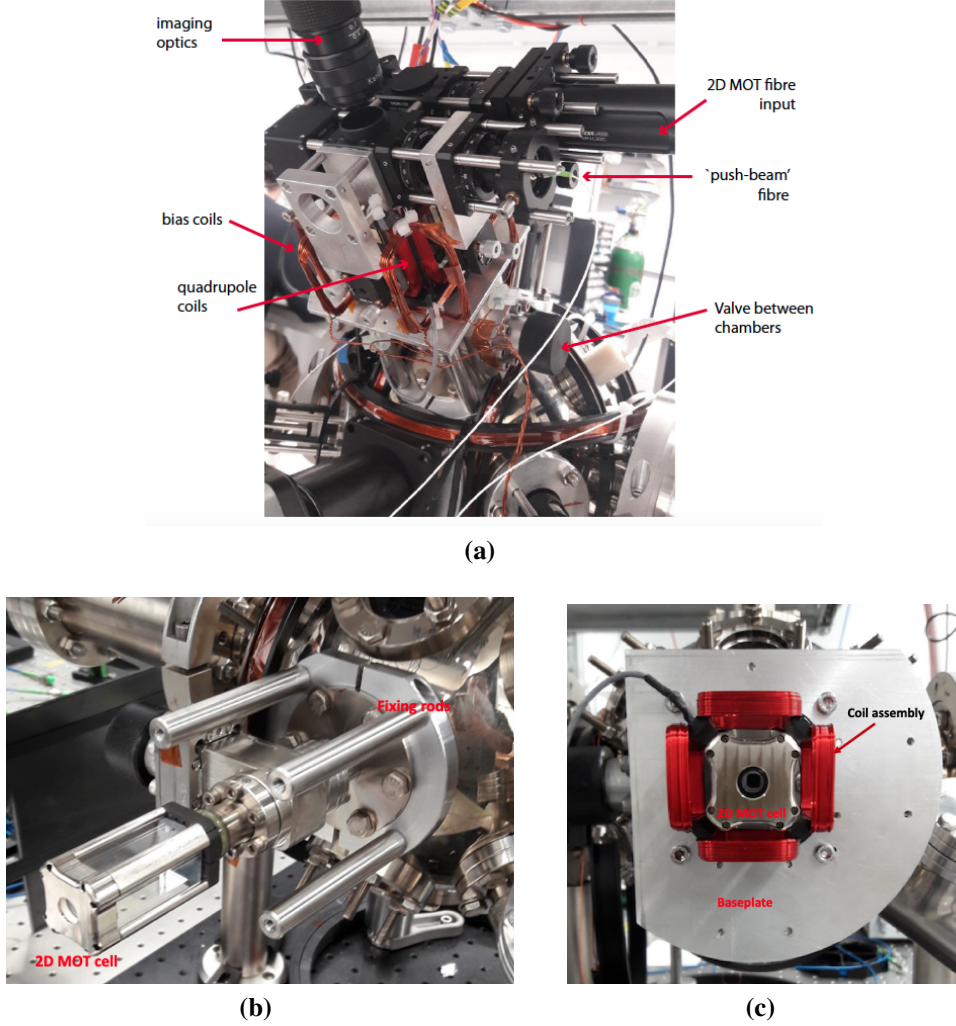


Figure 3.7: 2D-MOT assembly: (a) shows a picture of the entire 2D-MOT assembly, attached to the 3D-MOT interferometer chamber. Baseplate with optics and the coil assembly is mounted on a DN40 flange through four fixing rods (b). (c) is a picture taken along the long axis of the 2D-MOT with the baseplate and the coil assembly attached.

The 2D-MOT quadrupole coils and optics are designed and built tightly around the 2D-MOT glass cell on a separate baseplate with a diameter of 15 cm. The entire unit takes up a volume of less than 2 litres, as shown in Fig. 3.7 (a). The baseplate with optics and coil assembly is attached to one DN40CF flange of the interferometer chamber via four 12 cm long rods, as shown in Fig. 3.7 (b). (c) shows that the coils are free of contact with the glass cell. The length of the rods is chosen such that the optics form a MOT in the centre of the glass cell. Once cold atoms are formed, a push beam accelerates them into the 3D MOT interferometer chamber at an angle of 45° from the horizontal plane (X/Y -plane).

Quadrupole field and bias field coils

The quadrupole coil assembly of the 2D MOT is provided by *ColdQuanta Inc.* Two pairs of rectangular coils with inner diameters of 25.4 mm x 51 mm are orientated in a square structure with side length of 39 mm. This leaves enough space for the 2D MOT cell to slide through during installation, as shown in Fig. 3.7 (c). Opposite pairs of coils are wired up to produce opposite fields, giving a field gradient of 20 (G/cm)/A in the centre. Figure 3.8 illustrates the magnetic field lines (blue dashed lines) from a cross-sectional view. A fast IGBT switch controls the current going through the coils to ensure a short field decay time, because a slowly decaying residual field from the 2D MOT could reduce the efficiency of subsequent Sisyphus cooling in the main chamber by adding additional energy splitting between the magnetic-sensitive m_F states.

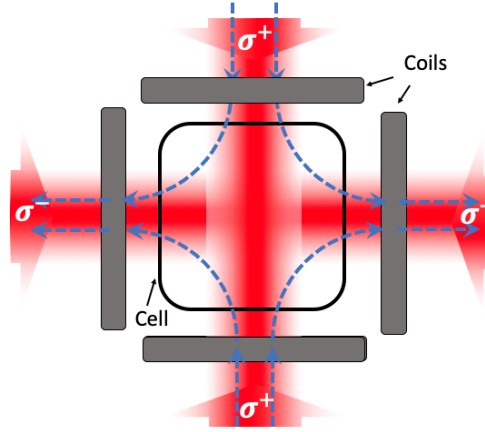


Figure 3.8: 2D-MOT cooling beams and magnetic field orientation.

Two additional pairs of bias coils allow us to adjust ambient field in the centre of the glass cell. These are made of 20 turns of 1 mm thick copper wires in rectangular shape (50 mm x 75 mm) and are separated by 55 mm from each other. Using these, we can vary the position of the 2D MOT to maximise 3D MOT loading through the aperture. A current change within ± 0.3 A shifts the 2D MOT across the aperture completely.

Optics

In the 2-dimensional magnetic quadrupole field, as illustrated in Fig. 3.8, opposite cooling beams are prepared in the same handedness, which drive the same σ^+ (or σ^-) atomic transition in the magnetic quadrupole field. The optics for the cooling beams are assembled using 30 mm cage components to distribute the light, as illustrated in Fig. 3.9 (a). (b) shows a schematic of the optics in the cage mount, in which I use XZ - and YZ -planes to illustrate the setup from different views.

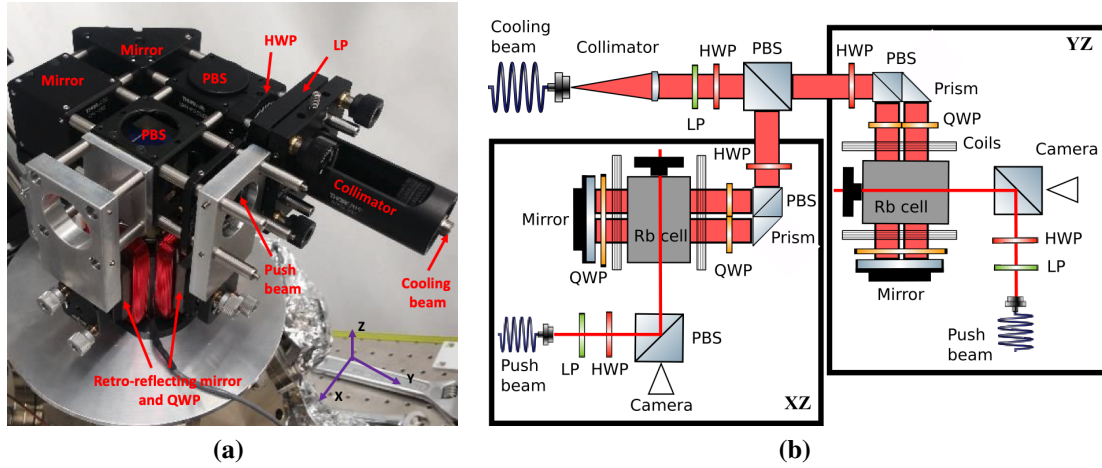


Figure 3.9: 2D MOT optics assembly shown in (a). A schematic of the setup is illustrated in (b).

The system takes two fibre inputs, one for the cooling/repump beam and the other one for the push beam. A single 1 inch lens collimates the diverging light from the cooling/repump fibre into a collimated Gaussian beam with a waist of 7.5 mm. A linear polariser (LP) purifies the polarisation before splitting the light into 2 beams with equal powers using a half wave plate (HWP) and a polarising beam splitter (PBS). The LP is essential to improve the stability of atom number. Any polarisation drift in the cooling/repump fibre induces unbalanced intensity fluctuations in the two split beams after the PBS. The LP transfers the unbalanced fluctuations into slow common mode intensity drift, which weakly affects the atom number in the 3D MOT.

These beams are directed, one along X and the other along Y . Each beam is then doubled to cover the full length of the cell along Z and divided by another HWP and a 15 mm PBS. The beam coming out of the reflected port of the 15-mm PBS goes through a quarter wave plate (QWP) into the 2D MOT cell. The beam transmitted by the 15-mm PBS is guided by a 15 mm prism via a second QWP into the cell. As the polarisations of these two beams are orthogonal, the two last QWPs need also to be aligned orthogonally in order to prepare the light in the same circular polarisation. To minimise any involvement of unnecessary components, the QWPs are cut and fixed in place so that no further alignment is required. On the opposite sides of the 2D MOT cell, mirrors (25 mm x 35 mm) with integrated QWPs retro-reflect these cooling beams.

In addition to the optics for the cooling, the light from the push beam fibre is collimated into a beam with a diameter of 1.5 mm. The collimator is mounted on a 1 inch kinematic mount with a 2-axis translation stage. Polarisation purification is necessary to stabilise the push beam power. A HWP and a PBS are used to deflect the push beam into the glass cell along its axis and to control the power. A camera on the other input port of the push beam PBS takes pictures of the 2D MOT, viewed along its long axis, for

diagnostics.

3.5.2 Characterisation

Following the specification of *ColdQuanta*, I set the quadrupole field of the 2D MOT to produce a magnetic field gradient of approximately 15 G/cm. The initial characterisation of the 2D MOT is carried out at a dispenser operating current of 2.5 A.

Frequency and power optimisation

The cooling frequency for the 2D MOT is optimised by maximising the number of atoms in the 2D MOT. Due to the densely packed optics, as shown in Fig. 3.9 (a), the only way to collect fluorescence photons from the 2D MOT is through the PBS along the Z -axis, which also guides the push beam into the cell. (See section 3.6.2 for further details of fluorescence detection.) A camera is setup, as illustrated in Fig. 3.9 (b), to capture images of the 2D MOT. As the images are taken along the longer axis of the 2D MOT, which exceeds the depth of focus of the camera, this measurement is not suited to accurately estimate the atom number. However, it captures very well the response of the atom number to the cooling frequency.

A 2D MOT is formed using the highest available optical power out of the 2D MOT fibre of approximately 70 mW. The camera is exposed for 5 ms for every image. To get an estimate of the atom number in the 2D MOT, I integrate all pixel counts within the region around the 2D MOT in the image and normalise the total counts against the scattering rate at the frequency detuning, which is used to form the 2D MOT. Figure 3.10 (a) shows this total pixel counts adjusted by the scattering rate as a function of the 2D MOT cooling frequency detuning from the resonance frequency of the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transition, which peaks between -11 and -14 MHz.

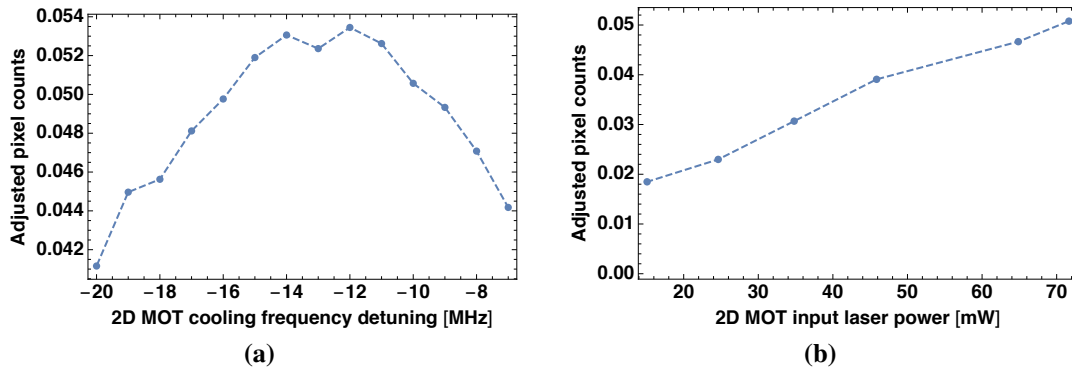


Figure 3.10: Frequency (a) and laser power (b) scans for the 2D MOT. The atom number in the 2D MOT is estimated from the total pixel count of every image taken from the 2D MOT, which is normalised against the scattering rate.

Next, after setting the cooling frequency detuning at -12 MHz, I vary the total input laser power in the 2D MOT and make the same proportionate measure of the atom number from the images taken at each input power. Figure 3.10 (b) shows a linear increase¹ in the adjusted pixel counts within the available laser power in the 2D MOT. This indicates that there is plenty room of improvement for the 2D MOT, if more laser power is available to us.

Push beam optimisation

The 3D MOT (see section 3.6) has a capture velocity of a few tens of m/s. Any atoms arriving at the 3D MOT with a velocity greater than that will not be captured. Because gravity acts on the atomic beam as it flies from the 2D MOT to the 3D MOT, atoms with a forward velocity of less than 6 m/s fall below the 3D MOT. (Here the capture volume of the 3D MOT is assumed to be limited by the size of the MOT beams.) Hence, to have good loading efficiency, it is important to maximise the number of atoms travelling between about 10 m/s and 40 m/s.

We have measured the time-of-flight distribution for atoms leaving the 2D MOT and arriving in the 3D MOT region, which is located 30 cm away. After loading for a time t_{load} of 10 ms, the 2D MOT is switched off. Next, I allow a flight time t before pulsing on all six 3D MOT beams, which are set on the resonance of the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transition for fluorescence detection using photodiode. The black line in Fig. 3.11 shows the number of atoms arriving in the detection region as a function of flight time. Atoms

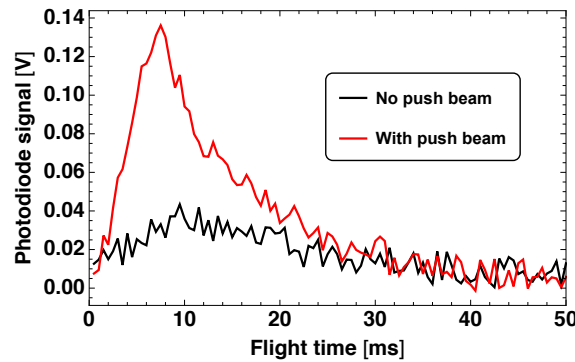


Figure 3.11: Time-of-flight distribution for atoms leaving the 2D MOT and arriving at the 3D MOT region. The 2D MOT is switched on for 10 ms together with (red) and without (black) the push beam.

travelling at desired velocity are expected to arrive at the 3D MOT region between approximately 5 ms and 30 ms. To increase the number of atoms arriving within this period of time, the push beam is switched on and off together with the 2D MOT. The red line

¹This exact linear trend is also observed, when looking at the atom number in the 3D MOT by varying the optical power in the 2D MOT.

in Fig. 3.11 shows the time-of-flight distribution for atoms leaving the 2D MOT with the push beam on. A clear increase in atom number is observed for atoms arriving earlier than 30 ms of flight time.

Next, we turn on the quadrupole coils of the 3D MOT so that atoms coming from the 2D MOT could be re-captured (see section 3.6 for further details of the 3D MOT). Here, the 2D MOT, push beam and the 3D MOT are turned on together for 50 ms of loading time, then we image the trapped atoms using fluorescence detection system in the main chamber (see section 3.6.2 for further discussion of the imaging system) and convert the photodiode voltage into number of atoms. We start with approximately 150 μW in the push beam and vary its frequency detuning from the $|F = 2\rangle \rightarrow |F' = 3\rangle$ transition. Figure 3.12 (a) shows the number of atoms trapped in the 3D MOT as a function of the push beam frequency. A clear maximum in the atom number in the 3D MOT is observed at +10 MHz. With the push beam frequency blue-detuned at approximately 10 MHz, I

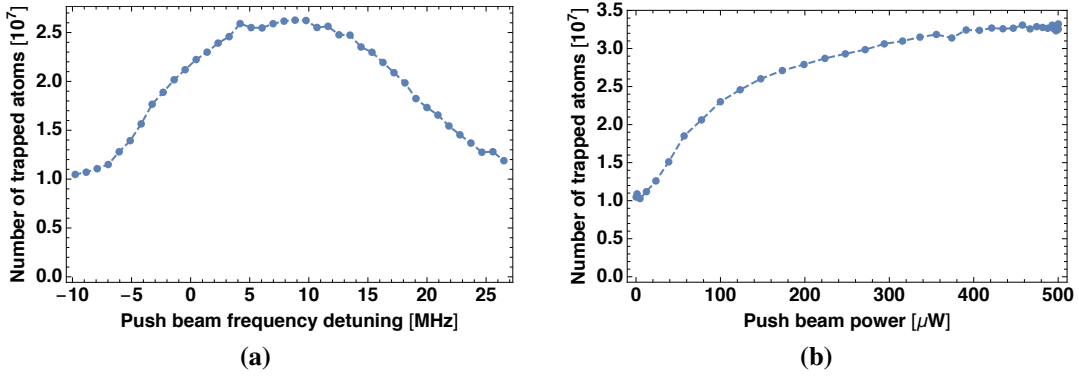


Figure 3.12: Optimisation of the push beam frequency (a) and power (b).

scan the number of trapped atoms as a function of the push beam power, as shown in Fig. 3.12 (b). This shows that the loading in the 3D MOT is saturated for the push beam power higher than 400 μW .

Dispenser current optimisation

Previous optimisation is carried out at a reasonably low Rb partial pressure in the 2D MOT cell, which is controlled by the current in the dispenser. Now, we change the dispenser current and turn on the 2D MOT and the 3D MOT together for a time t_{load} to monitor the number of trapped atoms in the 3D MOT. Figure 3.13 (a) shows the experimental data of the loading curves with the optimised push beam on (square) and off (circle), for several selected values of the dispenser current. In Fig. 3.13 (b), we summarise the loading rates $\gamma = 1/t_0$ as a function of the dispenser current, which are extracted from

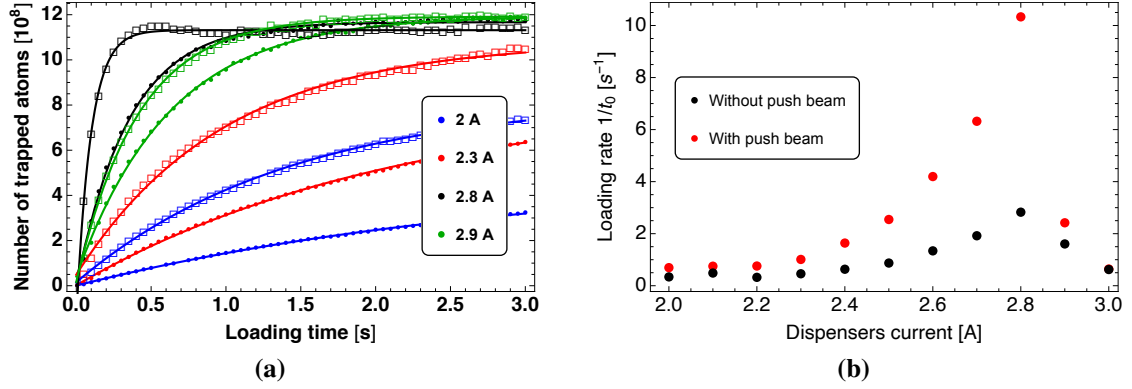


Figure 3.13: (a) shows the 3D MOT loading curves with the push beam switched on (square) and off (dots), respectively, as a function of dispenser current. (b) summarises the loading rate $1/t_0$ estimated from the fit (solid line) of the loading curves.

the fit in (a) (solid lines) using the following simplified equation [63]

$$N(t_{\text{load}}) = N_{\infty} (1 - e^{-t_{\text{load}}/t_0}), \quad (3.2)$$

where t_0 is the loading time. Black dots in Fig. 3.13 (b) show the loading rates due to diffusion through the aperture in the 2D MOT cell without the push beam, whereas the red dots indicate the loading with the push beam. Around 2.8 A, the loading rate peaks at 10 s^{-1} and 2.5 s^{-1} with the push beam on and off, respectively. Any further increase of the dispenser current rapidly decreases the loading rate. This occurs because the loss rate due to the high background pressure in the 2D MOT cell at a dispenser current higher than $> 2.8 \text{ A}$ outweighs the loading rate of the 2D MOT itself, which causes the 2D MOT to diminish. From these data, we can conclude that our 2D MOT is able to provide us with more than 10^8 atoms in the 3D MOT at a rate of 20 s^{-1} when using an input optical power to the 2D MOT of close to 70 mW (the maximum available to us at present.).

3.6 3D MOT

In the interferometer chamber, a main MOT is formed from the stream of 2D MOT atoms. This MOT has been referred to as the 3D MOT, as it is optically cooled on all three degrees of freedom. Analogously to the 2D MOT, cooling on each degree of freedom is achieved using a pair of counter-propagating beams [65], both equally red-detuned from the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ resonance. The negative frequency detuning is essential for laser cooling, as it compensates the Doppler shift of the transition frequency for atoms, which move towards the laser beams. To spatially confine the cooled atoms, a magnetic quadrupole field introduces a position dependence in the magnetic Zeeman splitting in

Rb. For counter-propagating laser beams with opposing polarisations ($\sigma^\pm$) or the same helicity (L or R), atoms in a given magnetic sub-level will be predominantly addressed by one of the beams, which drives the atom back to the centre of the trap. As the atom moves further away from the centre of the magnetic field, this restoring force on the drifting atom increases.

Figure 3.14 shows a schematic of a 3D MOT setup. A pair of electric coils are setup in anti-Helmholtz configuration, which provide a magnetic quadrupole field. Three pairs of the opposing laser beams are indicated with the thick arrows with corresponding polarisations. In our setup, the light for each cooling beam is prepared in the light distribution unit. In addition, we implement three separate pairs of bias coils around each axis to minimise the ambient magnetic field at the centre of the quadrupole field. This is important to perform efficient Sisyphus cooling.

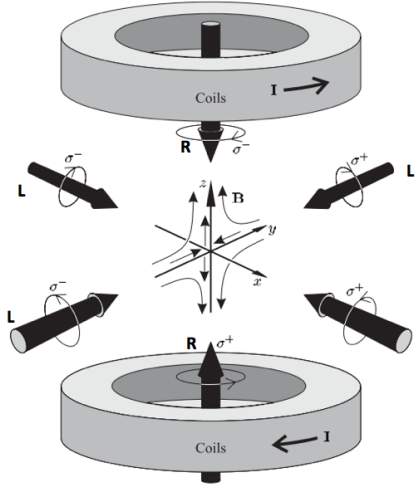


Figure 3.14: 3D MOT schematics: The handedness of each MOT beams are indicated by L and R respectively, whereas the polarisations of the light are labeled by σ^+ and σ^- . Figure taken from Ref. [65] and modified by the author.

Figure 3.15 summarises the sequential detail in the MOT forming stage. In order to load the 3D MOT the quickest, the magnetic field and laser light for the 3D MOT are switched on together with the field and the light for the 2D MOT and the push beam. Optimisation of the laser beams for the 2D MOT, the 3D MOT and push beam are carried out by maximising the atom number in the 3D MOT. The detuning of the cooling beams is set to -12 MHz from $|F = 2\rangle \rightarrow |F' = 3\rangle$ transition for the 2D MOT and -15 MHz for the 3D MOT with a magnetic field gradient of approximately 15 G/cm around the centre of the quadrupole field.

After filling the 3D MOT for 100 ms, the magnetic field and the laser light for the 2D MOT and push beam are turned off. We allow an additional 10 ms for the slow atoms to be captured in the 3D MOT. Then we switch off the quadrupole fields of the 3D MOT.

By reversing the voltage across the quadrupole coils during the switch-off, we are able to force the decay time of the field to be 0.8 ms. A total waiting time of 4 ms is added to the end of the MOT loading stage for the field to vanish completely, before moving on to the molasses phase.

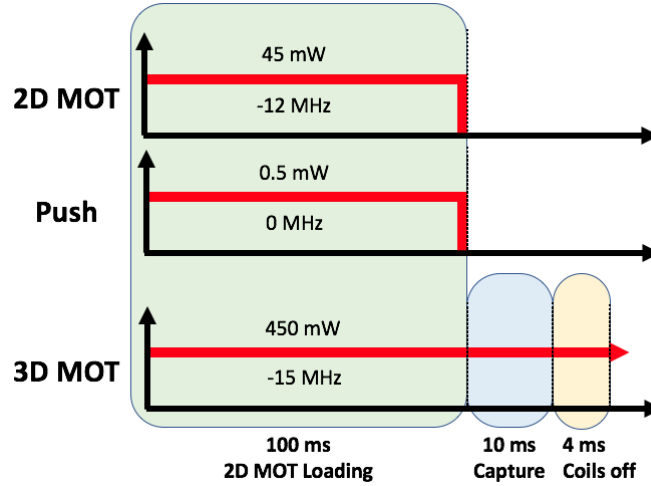


Figure 3.15: Experimental sequence during MOT loading stage for 2D and 3D MOT light.

In this section, I describe the necessary components to generate a 3D MOT in more detail and present a temperature measurement of the 3D MOT.

3.6.1 Setup

Quadrupole field coils

To improve the compactness of the setup and lower the power consumption, the quadrupole coils for the 3D MOT are designed to fit inside the vacuum chamber, which imposes some physical constraints. As the Raman beam optics, microwave view port and detection optics are all situated on the X/Y -plane in Fig. 3.4, the quadrupole coils are accommodated with their axis along Z (vertical) to avoid obstructing the other beam lines. The inner diameter of the coil formers should be at least 30 mm to leave room for the $Z_{\pm}$ MOT beams to propagate through. The outer diameter of the coil former is limited to 62 mm by the DN63CF flanges, which are the only openings to admit the in-vacuum coils into the interferometer chamber. Finally, the stream of 2D MOT atoms enters the 3D MOT chamber at an angle of 45° between the Z -axis and X/Y -plane. To avoid obstructing the 2D MOT stream, the separation between the coils must be at least 76 mm. Figure 3.16 (a) summarises all mentioned physical constraints. These leave radially 14 mm and axially 34 mm to distribute the coil wires. To achieve good heat conduction between the coils and the vacuum chamber, one end of the coil former is designed to match the spherical

inner surface of the vacuum chamber, as shown also in Fig. 3.16 (a). This enlarges the

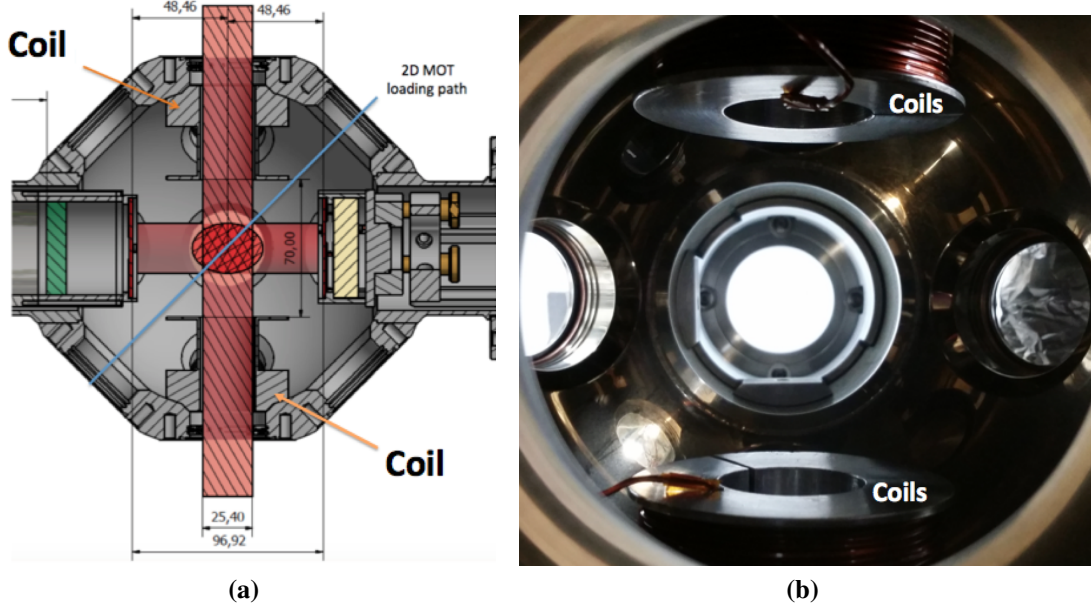


Figure 3.16: Schematics of the orientation of the MOT coils in the 3D MOT interferometer chamber (a). A picture (b) taken from the Raman axis shows the actual implementation.

contact area with the chamber and simplifies the mounting of the formers by attaching them to two DN40CF flanges using groove grabbers.

To minimise the electric power in the coils, we aim to reach a high packing density of the wires, which reduces the current for a given field gradient. For this purpose, UHV compatible copper wires with 1 mm square profile and 0.05 mm Polymide-ML lamination are purchased from *MWS Wire Industries*. In Fig. 3.17 (a), we use the Biot-Savart equation to calculate the dissipated electric power to achieve a field gradient of 15 G/cm in the centre, as required for a strong MOT. The calculations are carried out as a function of the coil separation Δz and number of axial turns, while fixing the number of radial turns at the maximum of 12. As we see, the power is highly dependent on the coil separation and is minimal for 20 axial turns. Guided by these simulations, we hand-wound two coils with 12 radial and 20 axial turns of wires around the coil formers, with a resistance of 0.5Ω each. Figure 3.17 (b) shows the magnetic field measured along the coil axis for one coil. The theoretical fit shown in red dashed line yields 2.71 ± 0.03 A of current, which is the only free fit parameter and agrees with the experimental setting of 2.6 ± 0.1 A within the error.

Figures 3.18 (a) and (b) show the measured magnetic field and the field gradient respectively, along a pair of quadrupole coils after being installed in the vacuum chamber. Without changing the current going through the coils, we fit the data for the magnetic field and the field gradient to independently determine the coil separation Δz . The estimated

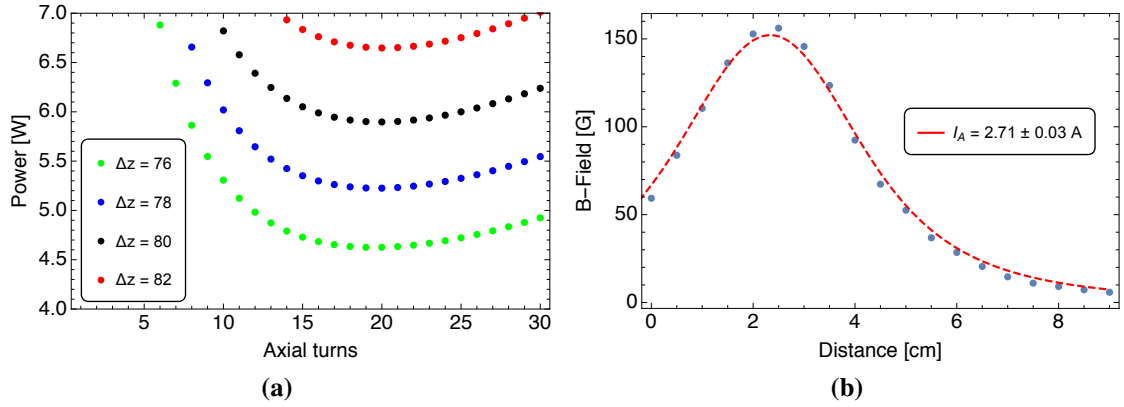


Figure 3.17: Simulation of power consumption (a) as a function of the coil separation and the number of axial turns of wires. (b) shows the measured magnetic field along the axis of a single coil, where red dashed line indicates the theoretical calculation with 2.7 A of current passing through.

separations from both fits agree with each other around 76 mm.

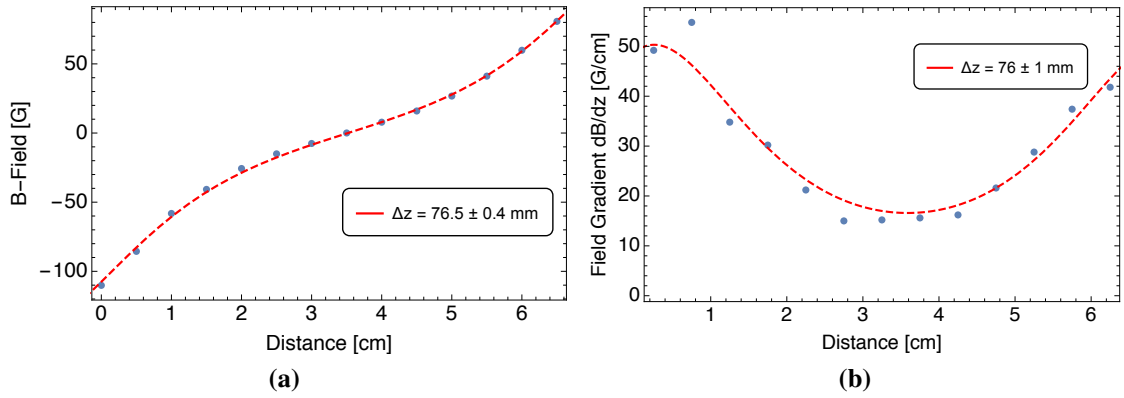


Figure 3.18: Fits (red dashed line) are applied to the data of the magnetic field (a) and field gradient (b) in Gauss to estimate the separation between the coils Δz . Both fits are in good agreement with each other.

If the heat produced by the coils can not be extracted from the vacuum, it could potentially melt the isolation layers of the coil wire and short the coils. To test this, we measured the temperature on the surface of the coils during continuous operation at 2.7 A, which is the operating current for the MOT. Figure 3.19 (a) shows the temperature rising from 22 °C room temperature to an equilibrium of 32 °C in red. By contrast, when the chamber is pumped out to a low pressure ($\approx 10^{-2}$ mbar), the temperature raises, but only to 39 °C. Similarly, Fig. 3.19 (b) shows double the cooling rate in air than in vacuum because of the additional effect of convection. These results validate the design of our quadrupole coils.

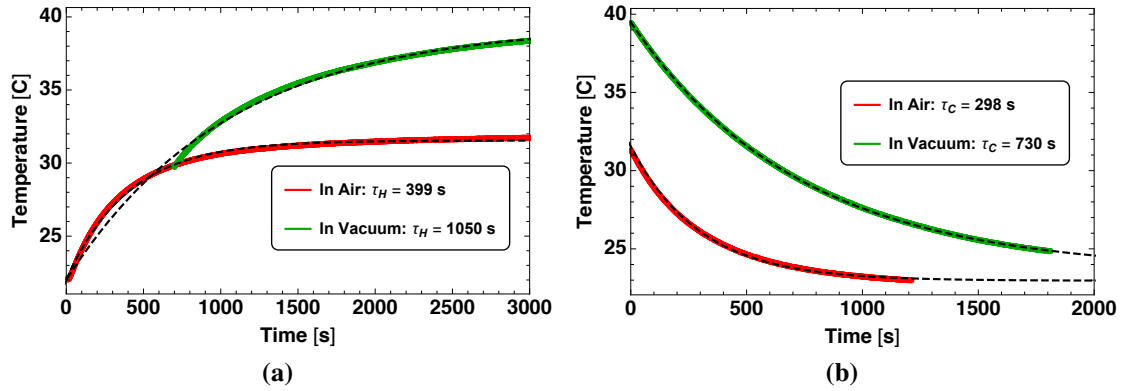


Figure 3.19: Temperature change of the coil during heating-up (a) and cooling-down (b) in air and in vacuum. The heating and cooling times (τ_H , τ_C) are obtained from exponential fits in dashed lines.

Bias field coils

Three pairs of bias coils in close Helmholtz configuration are designed to fit tightly on the 3D MOT interferometer vacuum chamber. Figure 3.20 shows an Illustrator drawing of the coils on the interferometer chamber without other components attached. These pairs of coils have different diameters in order to be fitted on top of each other. We choose the smallest coils with a diameter of 210 mm to be aligned along the Raman axis. Coils with 230 mm diameter are mounted vertically along the gravity (Z -axis). Finally, the largest coils with a diameter of 256 mm are oriented orthogonally to the previous two, along the axis of the detection optics. The formers for the bias coils are made out of PEEK to reduce weight. For each coil, 20 turns of 1 mm copper wires are hand-wound in-situ with the former attached to the chamber, giving a total resistance of 0.5Ω for the pair and a field of 1.5 G/A between the coils. As we use these bias coils to zero the ambient field

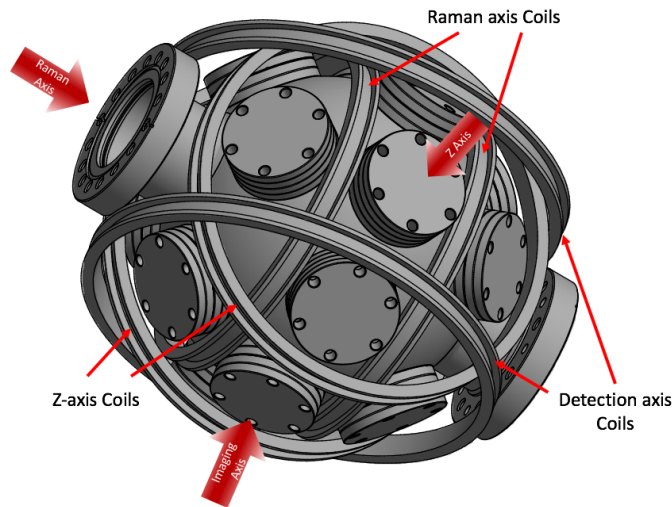


Figure 3.20: Schematics of the bias coils on the 3D MOT interferometer chamber.

during Sisyphus cooling and generate different quantisation fields for state preparation, microwave pumping, Raman interferometer and detection, field switching needs to be rapid, in order to minimise the measurement dead time. Therefore, three independent bi-polar current drivers are used to switch the field produced by each pair of bias coils with a rising time of 336 μs and a falling time of 142 μs respectively.

MOT beams

Six fibre outputs originate from the light distribution unit and provide us with the light for the 3D MOT. These beams have been referred as X_1 , X_2 , Y_1 , Y_2 , Z_+ and Z_- beams. Four separate AOMs are used to gain independent control over their frequencies and powers, as described in section 3.4. This flexibility allows us to launch, to detect atoms and to perform state preparation (see section 3.8). During the MOT loading stage, the frequencies and powers of all six beams are made equal in order to balance the scattering forces of each beam on the atoms. During the launch (see section 3.7.1), the frequencies of counter-propagating beams are oppositely detuned. During the state preparation and detection, only the Z -axis beams are used, while the Z_- -beam alone is employed for state clearance (blow-away).

Similarly to the collimator in the 2D MOT optics assembly, we use 30 mm cage components to build 3D MOT collimators. A fibre plate is mounted on a rotational cage mount with a linear polariser, which is aligned to the polarisation of the light coming out of the fibre and improves the polarisation purity. A single 7.5 mm convex lens collimates the beam into a zero-order quarter wave plate (QWP). The collimation is carried out using a Sheer-plate interferometer. Figure 3.21 shows the waist estimation for all six collimators from images taken at 15 cm away from the QWP, which corresponds to the position of the atoms in the chamber. All collimators are combined with kinematic mounts for more adjustment and attached to an adapter plate on the DN40CF viewports on each axis.

The polarisation of each collimator can be adjusted by rotating the fibre plate together with the linear polariser relative to the QWP. All collimators on the X/Y -plane are set to give out left-handed circularly polarised light, while the collimators on the Z -axis output right-handed circularly polarised light. We quantify the purity of the circularly polarised light coming out of each collimator by measuring its projection onto linear polarisation. This is verified by passing the collimated beam through a high-quality polarising beam splitter (PBS) with a polarisation extinction ratio (PER) of ≤ -30 dB. For ideal circularly polarised input light before PBS, the transmitted optical power should not oscillate as a function of PBS angle, Θ . Any impure, circular polarisations in the output beam are either introduced by a misalignment angle, α , between the optical axes of the QWP and

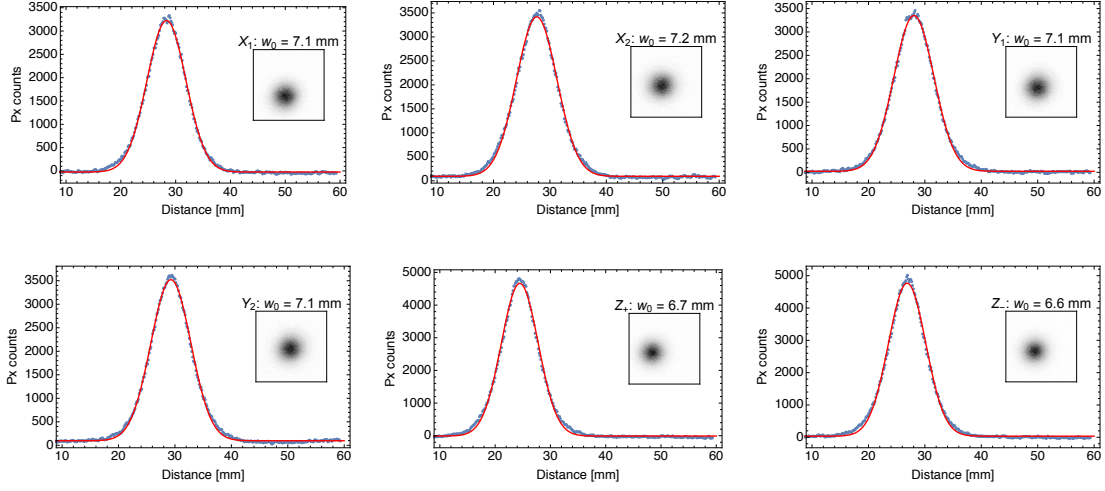


Figure 3.21: Beam waists measurement of the MOT collimators

the linear polarisation in the input beam, or caused by ellipticity in the input beam.

Here, we denote the laser light coming out of the fibre as $|I_0\rangle$. We describe $|I_0\rangle$ in a generic form as

$$|I_0\rangle = \frac{1}{\sqrt{a^2 + b^2}} (a |H\rangle + b e^{i\phi} |V\rangle), \quad (3.3)$$

where $|H\rangle$ and $|V\rangle$ are two orthogonal fibre axes. a and b are the respective electric field amplitudes. ϕ is the phase factor between the two linear polarisations. The light propagates through a QWP and a PBS. The transmitted light intensity I can be written as

$$I(\alpha, \Theta) = |\langle H | R(\Theta) \cdot Q \cdot R(\alpha) | I_0 \rangle|^2, \quad (3.4)$$

where $R(\alpha)$ is the rotation matrix addressing the misalignment between the QWP and the fibre axis. Q is the Jones' matrix for a QWP and $\langle H | R(\Theta)$ describes the orientation of the PBS. Equation 3.4 shows that the transmitted intensity through the PBS oscillates, as we vary Θ . Therefore, we use visibility $V = (I_{\max} - I_{\min}) / (I_{\max} + I_{\min})$ to relate this oscillation to the intensity ratio of wanted and unwanted polarisations, $r = b^2/a^2 = I_V/I_H$. Figure 3.22 exhibits the calculated visibility as a function of the intensity ratio r . For perfect QWP alignment, when $\Theta = 45^\circ$, Eq. 3.4 is reduced to the following simpler form

$$I(\Theta) = \frac{1}{2} + \frac{ab \sin(2\Theta + \phi)}{a^2 + b^2}, \quad (3.5)$$

where visibility can be written as $V = 2\sqrt{r}/(1+r)$. This indicates that visibility vanishes only for linearly polarised light coupled in the horizontal fibre axis and, most importantly, visibility becomes independent of the phase ϕ , as shown in the black line in Fig. 3.22. For a misaligned QWP by 2° , visibility reaches the minimum at an intensity ratio of

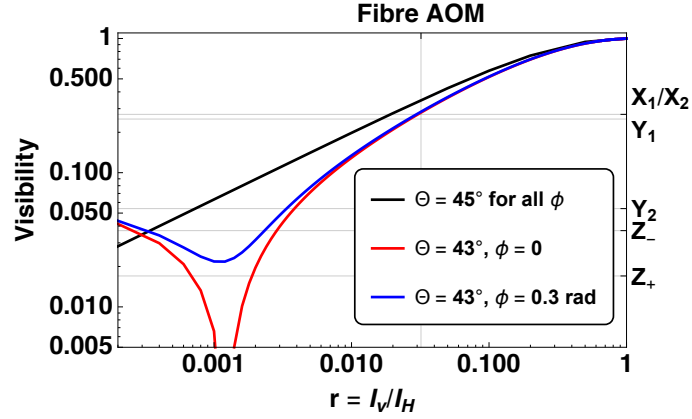


Figure 3.22: The visibility in Eq. 3.4 is plotted as a function of the intensity ratio r . The black line plots the case of perfect QWP alignment, where visibility is independent of ϕ . Red and blue show the visibility with a 2° misalignment between the QWP and the fibre axis, where the presence of ϕ increases the visibility additionally. The vertical gridline sets the highest expected visibility due to the quality of the fibre AOM and the horizontal lines indicate the measured visibilities after all MOT collimators.

approximately 0.001. This means that the misalignment can be compensated by a rotation in the input beam with a pure linear polarisation, as shown in the red line ($\phi = 0$). For non-vanishing ϕ , when ellipticity exists in the input beam, visibility increases for a fixed intensity ratio r , as illustrated in the blue line.

To optimise the polarisation of the MOT beams after the QWP, we minimise the visibility of the transmitted optical power through the measurement PBS by rotating the QWP. Simultaneously, to change ϕ , we heat up the fibre, which guides the light into the MOT collimator. An optimal alignment of the QWP is reached, when the visibility is also least sensitive to the temperature change of the fibre. As the fibre AOMs² have the highest polarisation extinction ratios in the light distribution, we expect that the purity of the polarisation after each MOT collimator is limited by the AOMs. Using the specification of the fibre AOMs, we expect an intensity ratio of unwanted and wanted polarisations to be 0.03, which is indicated by the vertical gridline in Fig. 3.22. This means that for perfect QWP alignment, the highest visibility, measured after each collimator, is expected to be around 0.35. Fortunately, the visibilities measured at all six collimators are lower than the expectations, as shown in Fig. 3.22 by the horizontal gridlines.

3.6.2 Imaging systems

It is convenient to use fluorescence detection to monitor the atom cloud. We measure the size of the cloud using a CCD camera and the number of atoms using a low-noise photodiode. The collection optics are outside the vacuum chamber, but close to a vacuum window in order to subtend the most available solid angle. This is achieved by using

²Our fibre AOMs are specified by a PER of ≤ -15 dB.

60 mm cage system mounted on an adapter plate on a DN40CF viewport on the interferometer chamber. A 2-inch convex lens ($f = 150$ mm) collects about 0.7 % of the light scattered from the atoms and collimates it into a 2-inch convex lens ($f = 75$ mm). A third 2-inch lens with a focal length of 60 mm is placed at 10 mm away from the 75 mm lens. Using ABCD matrices, we calculate that the system forms an image with a magnification of $-1/4$. This demagnification is necessary to provide enough field of view, while keeping the image within the detection area of the photodiode.

Acquiring the size of the atom cloud

A CCD camera is particularly important when performing temperature measurement, where the spatial information of the atom cloud is required. A CCD camera (*AVT PIKE F-505B*) with 2452×2054 pixels is mounted on a 3-axis translation stage at the output of the imaging system to captures the expanding size of the cloud as a function of time. The acquisition software defines the exposure time of the CCD camera by the duration of an external TTL signal. We have checked the distance in the object plane corresponding to one pixel on the CCD by taking a series of images of the MOT falling under gravity. From the time-of-flight positions of the centre of mass of the MOT, we have determined that one pixel corresponds 13.59 ± 0.05 μm in the object plane. The manufacturer gives a pixel size of 3.45 μm , implying a magnification of $-1/3.94$, which is very close to the design. This calibration is continuously updated a priori every temperature measurement of the MOT to assure the alignment of the imaging system.

Measuring the atom number

A low-noise photodiode, *FEMTO* LCA-S-400K-SI, with a detector area of 2.5 mm in diameter is mounted at the place of the camera to measure the number of atoms. We estimate the number of atoms N_{atom} by measuring the optical power of the fluorescence photons from the atom cloud. Fluorescence is driven by two frequencies: a strong imaging frequency and a weak repump frequency. The imaging frequency is set close to the $|F = 2\rangle \rightarrow |F' = 3\rangle$ transition, where the fine detuning from the resonance is described by angular frequency δ . The repump frequency is resonant to the $|F = 1\rangle \rightarrow |F' = 2\rangle$ transition, which pumps ($F = 1$) atoms into ($F = 2$) for detection. The scattering rate R_{sc} of an atom is

$$R_{\text{sc}} = \frac{\Gamma}{2} \cdot \frac{s}{1 + s + (2\delta/\Gamma)^2}, \quad (3.6)$$

where Γ is the $2\pi \times 6.07$ MHz natural line width of the ^{87}Rb D2 line and s is the intensity of the imaging beams in the units of saturation intensity. To converting the photodiode voltage V_{PD} into optical power, we use the photodiode module gain G_{PD} (5.3×10^6 V/W

at 780 nm), the optical acceptance solid angle, Ω , and the transmission efficiency, η_0 , of the imaging optics. These are all summarised in Eq. 3.7 to describe the number of atoms:

$$N_{\text{atom}} = \frac{4\pi}{\hbar\omega R_{\text{sc}}(\delta, s) \Omega \eta_0} \cdot \frac{V_{\text{PD}}}{G_{\text{PD}}}, \quad (3.7)$$

where $\hbar\omega$ is the energy of a scattered photon. Using Eq. 3.7, we have determined that our 3D MOT can load at least 10^8 atoms in less than 100 ms (see section 3.5.2), providing enough atoms for interferometry.

3.6.3 Characterisation

Temperature measurement

The temperature of the MOT is measured by ballistic expansion. As the MOT expands after the magnetic and optical fields are switched off, we record the fluorescence images of the cloud as a function of time t and measure the width $\sigma(t)$ of the cloud, which we take as the standard deviation of the Gaussian intensity distribution. From this, we deduce the cloud temperature T using

$$\sigma^2(t) = \sigma_0^2 + \frac{k_B T}{m_{\text{Rb}}} t^2, \quad (3.8)$$

where σ_0 is the initial cloud width, k_B is the Boltzmann constant and m_{Rb} is the atomic mass of ^{87}Rb .

The velocity distribution is not always isotropic. To determine the “temperature” along the x (horizontal) direction, we integrate over the image of the cloud in the y (vertical) direction. We then obtain the width by fitting the integrated signal to the function

$$c(x) = c_{\text{off}} + c_0 e^{-\frac{(x-x_0)^2}{2\sigma^2}}, \quad (3.9)$$

where $c(x)$ is the pixel counts at the calibrated position x on the CCD chip, c_{off} is the background pixel counts, c_0 is the amplitude and x_0 is the displacement of the centre of the mass of the cloud. Conversely, we integrate over x to find the width along y . Figure 3.23 shows the linear fits of Eq. 3.8 to the widths of the MOT on both axes, which yield temperatures of:

$$\begin{aligned} T_x &= 230 \pm 2 \mu\text{K} \\ T_y &= 371 \pm 3 \mu\text{K}. \end{aligned} \quad (3.10)$$

These velocity distributions are wider than we need for operating the interferometer and

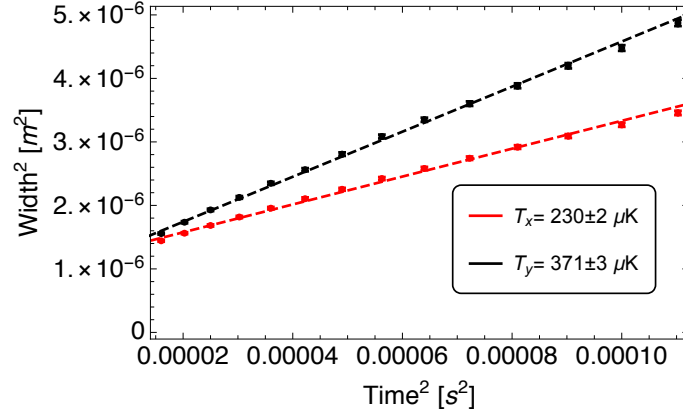


Figure 3.23: Temperature measurement of a MOT. Each data point represents the spatial width of the cloud, estimated from an image exposed for 0.5 ms by driving $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ and $|F = 1\rangle \leftrightarrow |F' = 2\rangle$ transitions. The error bars are estimated using the largest contribution, which is the uncertainty during the CCD chip calibration.

therefore we employ a second phase of cooling using optical molasses.

3.7 Molasses

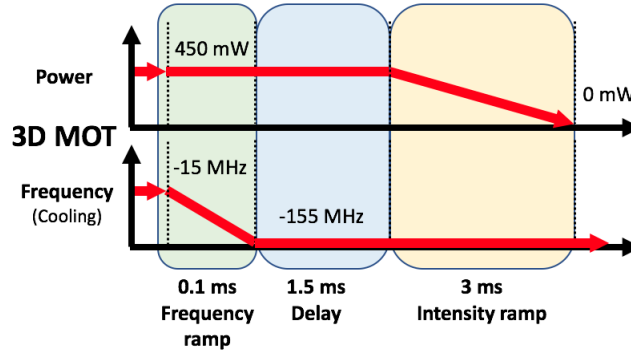


Figure 3.24: Experimental sequence for sub-Doppler cooling.

Optical molasses uses Sisyphus cooling, or polarisation gradient cooling is able to bring the temperature of the cloud down. Figure 3.24 summarises our optimised experimental settings to achieve a $8 \mu\text{K}$ cloud. At the end of the MOT loading phase, the 3D quadrupole coils are switched off. A delay time of 4 ms is added to the sequence for the residual magnetic field to decay away completely. As the final molasses temperature is proportional to $I/|\delta|$ [31], we increase the red detuning δ and then ramp down the intensity I of the laser in order to lower the temperature.

During the frequency ramp, we have empirically optimised the sweep rate and duration by minimising the molasses temperature. Here, the detuning of the cooling light frequency is swept from -15 MHz to -155 MHz over 0.1 ms. At the same time, the

repump frequency is kept on resonance on the $|F = 1\rangle \rightarrow |F' = 2'\rangle$ transition throughout to ensure that atoms do not fall into a dark state. It is important to mention that the repump light is generated by phase-modulating the cooling light using an EOM, as described in section 3.2.1. This phase modulation introduces sidebands at 6.8 GHz relative to the cooling frequency. The power in the repump beam depends on the modulation depth of the EOM, which, in turn, depends on the modulation frequency, as well as the cooling light frequency itself. Therefore, in order to ensure that the power of the cooling and repump beams remain constant during the frequency ramp, we also ramp down the modulation depth of the EOM. Following the frequency ramp, an additional delay time of 1.5 ms allows the serial commands for the frequency ramps to be completed by the laser.

After the frequency ramp has finished, we simultaneously reduce the intensities of all MOT beams over 3 ms. We have found that a linear intensity ramp in time minimises the atom loss. This is achieved by ramping the rf amplitude, which goes into the output AOM in the μ Quans laser, according to a measured calibration table between the rf control voltage and the output laser intensity. Figure 3.25 shows a screenshot of oscilloscope traces for the laser intensity (green) and the rf control voltage (blue). After this linear intensity ramp, we have found that approximately 90% of atoms prepared in the MOT still remain in the trap and ready for the state preparation.

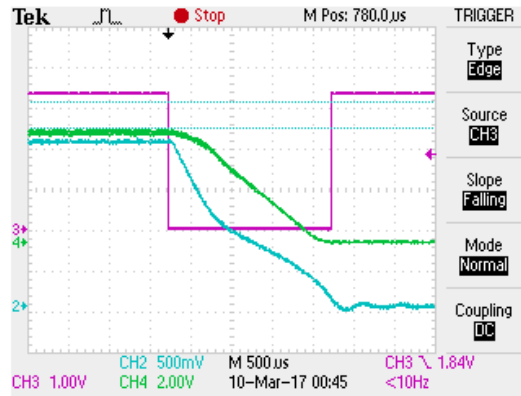


Figure 3.25: Change in the laser power (green) with the rf control voltage (blue). The purple line is the trigger for the scope.

3.7.1 Launching

The necessity for launching atoms arises from the way we arrange the laser beams that drive the interferometer Raman transitions. Figure 3.26 (left) illustrates the four beams seen by the atoms. Beams 1 and 4 are the incoming Raman beams and 2 and 3 are the reflected beams (further details in chapter 4). Red and blue represent their frequencies.

These four beams allow four separate Raman transitions. The beams 1,3 and 2,4 are both counter-propagating and therefore the transitions they drive exhibit a Doppler shift if the atoms have a velocity component along the beams. These transitions will be referred as velocity-sensitive. However, the co-propagating beam pairs 1,4 and 2,3 drive Raman transitions with no significant Doppler shift. These will be referred as velocity-insensitive transitions. Figure 3.26 (right) shows a Raman spectrum for a cloud of cold atoms moving to the right in Figure 3.26. The large peak at the centre is the velocity-insensitive transition caused by the co-propagating beams. The broad peak on the left is the Doppler-shifted resonance driven by beams 1 and 3, while the peak on the right is driven by beams 2 and 4.

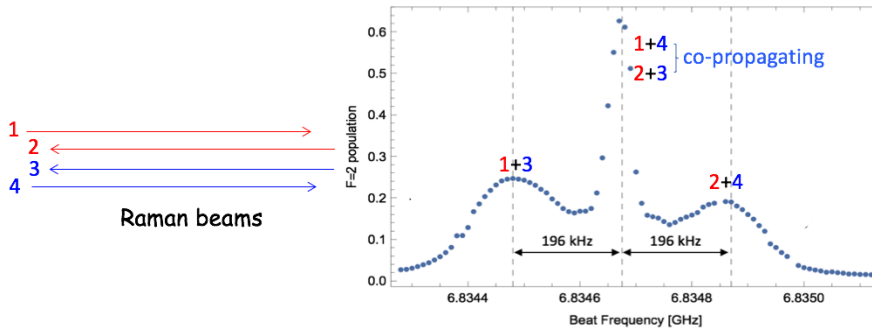


Figure 3.26: Orientation of the Raman beam in retro-reflection configuration (left). Raman spectrum scan (right) by scanning the beat frequencies between the two Raman lasers.

To make the interferometer work properly, we need to drive the Raman transitions with only one beam pair (either 1,3 or 2,4). We achieve this by launching the cloud along the Raman beam axis using the moving molasses technique, and tuning the Raman beat note, the effective Raman laser frequency, up or down (by 196 kHz in this example) to be resonant with atoms at that velocity.

To make the molasses move along one of the molasses (MOT) beams, we detune the frequency of that beam by $-\delta f$ (red-detuned), while the counter-propagating beam is detuned by $+\delta f$ (blue-detuned). The atoms are then cooled to rest in the frame moving with velocity $\delta v = \lambda \delta f$. To launch along the Raman axis, we detune both pairs of horizontal molasses beams by $\pm \delta f_R$ giving a launch velocity of

$$\delta v = \sqrt{2} \lambda \delta f_R. \quad (3.11)$$

This is illustrated in Fig. 3.27, where we have also included a vertical component of launch velocity controlled by δf_z . Because of the acceleration under gravity, the vertical launch helps to keep the atoms within the Raman beam region for a longer time. In addition, the vertical launch velocity is set such that the atom cloud comes back to its

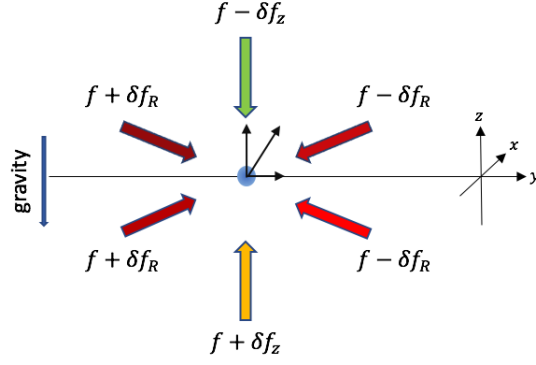


Figure 3.27: Orientations of the MOT beams and their frequency detunings for launching. The atoms move on the yz -plane.

initial position after around 50 ms of flight time. This simplifies the settings of $\pi/2$ - and π -pulses of the interferometer, see section 5.1.

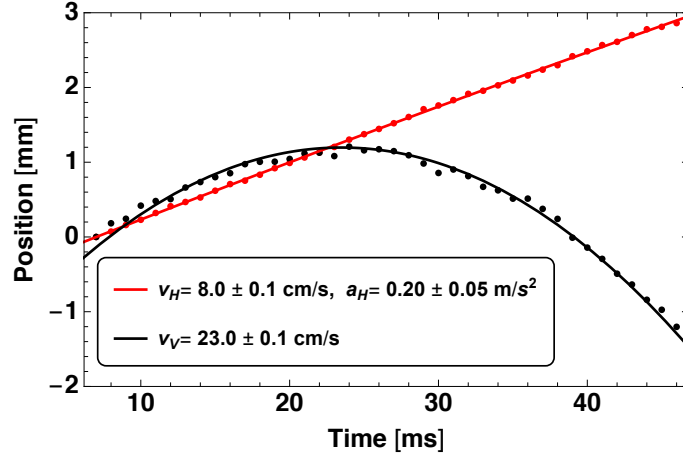


Figure 3.28: Horizontal and vertical positions of the atoms after launching.

We use a detuning of $\delta f_R = 75$ kHz, which corresponds to a horizontal launch velocity of $v_H^{(0)} = 8.3$ cm/s along the Raman axis. The detuning along the Z -axis is set to $\delta f_z = 320$ kHz, corresponding to a vertical velocity of $v_V^{(0)} = 25.0$ cm/s. The dots in Fig. 3.28 show the horizontal (red) and vertical (black) positions of the launched atom cloud, measured as a function of time using the CCD camera. These are fitted to the functions:

$$\begin{aligned} x_H(t) &= x_{H,0} + v_H t + a_H \frac{t^2}{2}, \\ x_V(t) &= x_{V,0} + v_V t - g \frac{t^2}{2}, \end{aligned} \quad (3.12)$$

where v_H and v_V are the initial launch velocities of the cloud motion. A horizontal component of acceleration a_H is included to $x_H(t)$, in order to account for a possible

small misalignment angle between the camera and the gravity $\vec{g}$. The fit gives $v_H = 8.0 \pm 0.1 \text{ cm/s}$ and $v_V = 23.0 \pm 0.1 \text{ cm/s}$. These differ from the ideal 8.3 cm/s and 25.0 cm/s , perhaps because of a residual magnetic field δB , which would shift the final velocity of the molasses by $\mu_B g_F \delta B / \hbar k$ [66], as the Doppler shift needs to compensate for the Zeeman shift. Stray magnetic fields of only 3 mG^3 and 50 mG respectively along the Raman and Z -axis would explain the observed deviations. This seems very likely. The fit in Fig. 3.28 shows a horizontal acceleration of $a_H = 0.20 \pm 0.05 \text{ m/s}^2$, which we think is due to a misalignment angle of around 1° .

3.7.2 Molasses temperature measurement

To ensure the coldest molasses temperature, it is important that the magnetic field is as small as possible at the location of the atoms. We therefore also null any ambient laboratory fields using the three pairs of Helmholtz bias coils shown in Fig 3.20. By measuring the temperature of the molasses for different bias settings, we can then ensure the coldest cloud possible. Figure 3.29 shows the expansion of a cold cloud falling under gravity along three different trajectories: without launching (top), with horizontal launching along the Raman axis (middle) and with horizontal and vertical launching along the Raman axis and against the gravity (bottom) respectively. As described in section 3.6.3, we

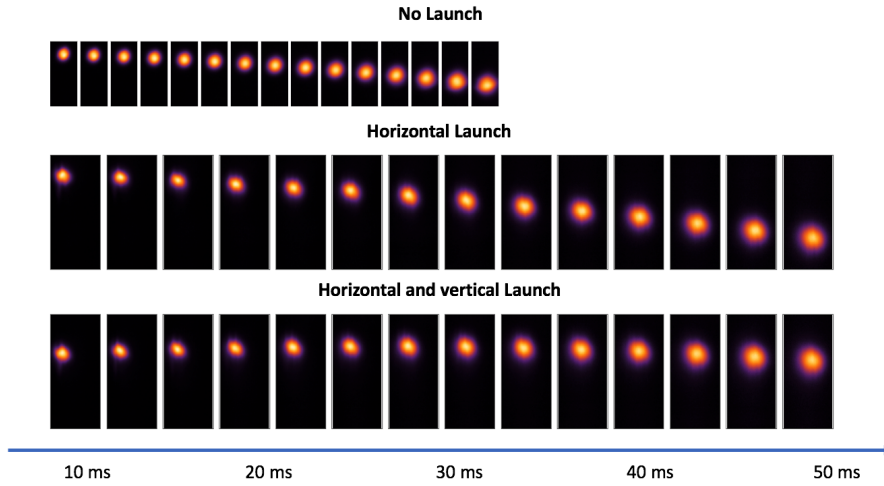


Figure 3.29: Expansion of molasses for times between 10 ms and 50 ms. Sequential images are taken without applying any launching (top), with horizontal launching only (middle) and with horizontal and vertical launching (bottom).

can determine the temperature of the atoms by measuring the expansion rate of this cloud after the laser light and the magnetic field are switched off. Figure 3.30 shows the widths of the cloud at different expansion times for all three launch settings. From the rates of

$^3 1 \text{ G} = 10^{-4} \text{ T}$

expansion, we determine that the cloud has a temperature of around $6-8\ \mu\text{K}$ on both axes in all three cases. A small increase in the temperature was observed with launching applied. This, however, would insignificantly affect the number of velocity-selected atoms in the interferometry, when the selection pulse is applied on resonance (see Chapter 5).

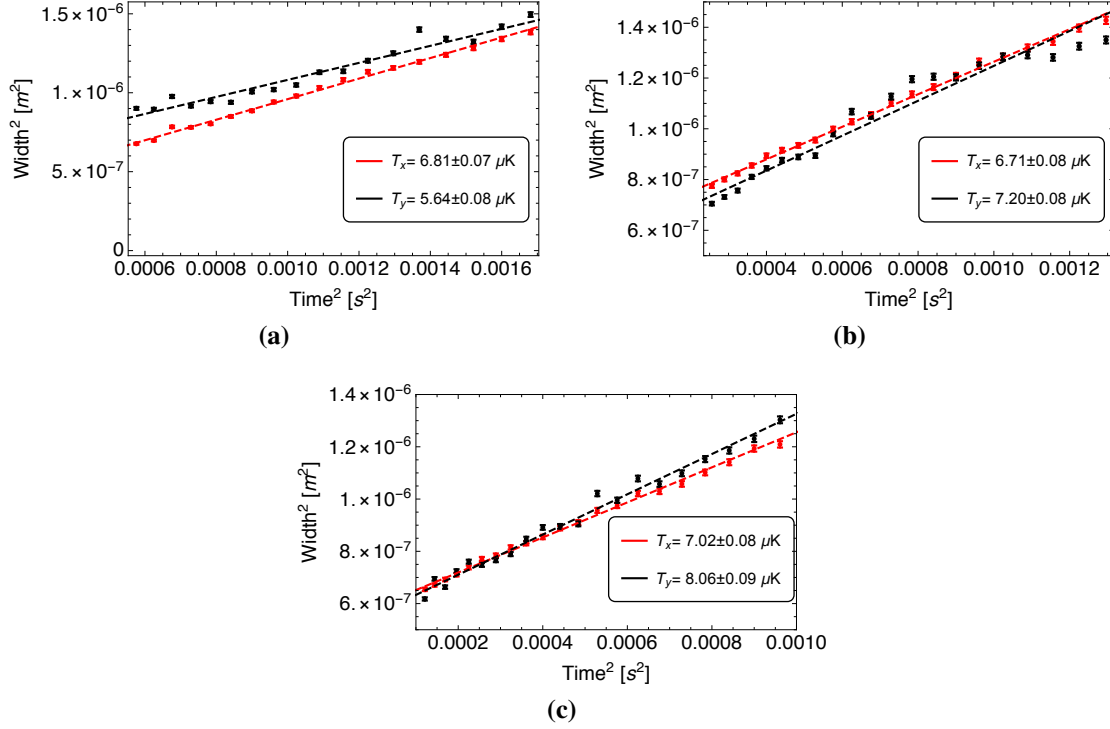


Figure 3.30: Expansion of a cold cloud after molasses in three launching configurations: no launching (a), horizontal launching (b) and horizontal and vertical launching (c).

3.8 State preparation

To make the interferometry robust against ambient magnetic field fluctuation, interferometer pulses act on the magnetically insensitive hyperfine Zeeman sub-levels ($F = 1, m_F = 0$) and ($F = 2, m_F = 0$) in the ground state of ^{87}Rb . At the end of the polarisation gradient cooling, the majority of atoms are distributed between the five Zeeman sub-levels of $F = 2$ hyperfine level, with the rest on the sub-levels of $F = 1$ hyperfine level. To increase the number of atoms in the magnetically insensitive states, we apply a sequence of steps to concentrate all population into the ($F = 1, m_F = 0$) state, before selecting a narrow velocity class of atoms for the interferometry. These steps are illustrated in Fig. 3.31.

In summary, four steps are used to achieve a high population in the ($F = 1, m_F = 0$) state. Step 1 and step 2 are purely optical. In the step 1, all $F = 2$ atoms are optically

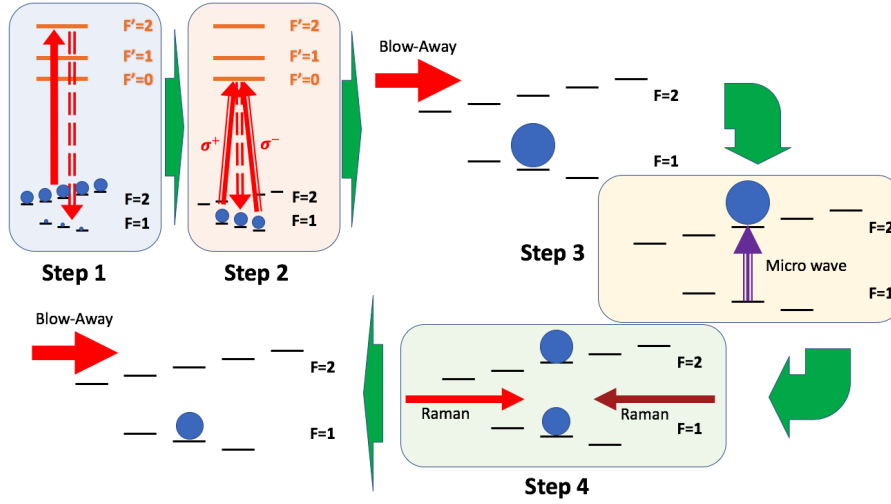


Figure 3.31: Schematics of the 4-step state preparation. Step 1 and 2 optically drive the population in the molasses into $(F = 1, m_F = 0)$ state. After blowing away residual atoms on $F = 2$, a microwave pulse transfers all population to $(F = 2, m_F = 0)$ state in the step 3, which are then velocity selected back into $(F = 1, m_F = 0)$ by a long Raman pulse in the step 4. A second blow-away beam purifies the population for the interferometry.

pumped via $F' = 2$ into the three magnetic sub-levels of $F = 1$ state. In the step 2, we use a pair of counter-propagating σ^+/σ^- beams to pump the population via $F' = 0$ into the dark state $(F = 1, m_F = 0)$. A blow-away beam tuned to the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transition then clears out any population remaining in the $F = 2$. Step 3 is a microwave transition, which transfers all atoms in $(F = 1, m_F = 0)$ to $(F = 2, m_F = 0)$. A long velocity-selective Raman pulse (see Chapter 5) then transfers a narrow velocity class of atoms from $F = 2$ to $F = 1$ for interferometry in step 4. The remaining $F = 2$ atoms are again cleared out by a second blow away beam. Figure 3.32 shows the details of the 4-step state preparation sequence, starting from the end of the polarisation gradient cooling. We begin with a 4 ms delay, which allows a quantisation field to be established along the optical axis of the counter-propagating Z -beams for use during step 1 and 2. Then, the quantisation field is adiabatically rotated to align with the magnetic field of the microwave so that it drives the $\Delta m = 0$ transition efficiently. After a $80 \mu\text{s}$ microwave pulse, we allow another 4 ms for the B field to be established along the Raman axis for use in step 4. In the sections 3.8.2-3.8.4, I provide further details about each of these steps, but first, I describe our state-sensitive detection method.

3.8.1 State-sensitive detection

As described earlier in section 3.6.2, to perform state selective detection on atoms, we induce fluorescence from the $F = 2$ atoms, with the atom number N_2 . This is achieved by a short pulse, which resonantly drives the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transition. The photodiode

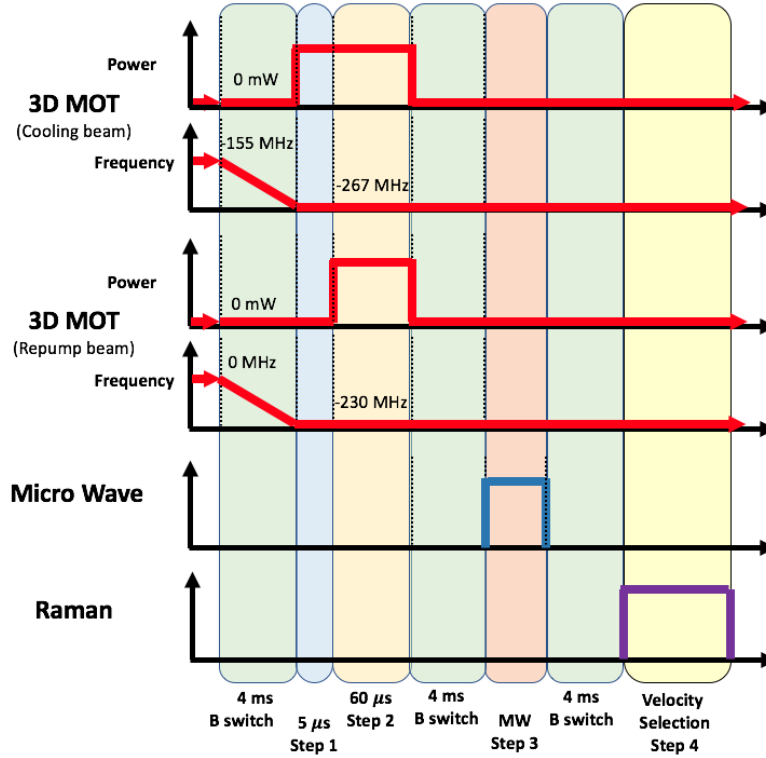


Figure 3.32: Experimental sequence for the 4-step state preparation.

converts the measured optical power into voltage V_2 . We also measure the total number of atoms by simultaneously driving the $|F = 1\rangle \leftrightarrow |F' = 2\rangle$ and the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transitions, which recycles $F = 1$ atoms also to $F = 2$ for detection. The total number of atoms, N_{tot} , is measured after the N_2 measurement. Next, the atoms are all blown away by a resonant beam, and then, the background signals, B_2 and B_{tot} , are collected with the same conditions⁴ for the detection beams as were used to measure N_2 and N_{tot} . This detection sequence is illustrated in Fig. 3.33. With this scheme, we are able to derive the

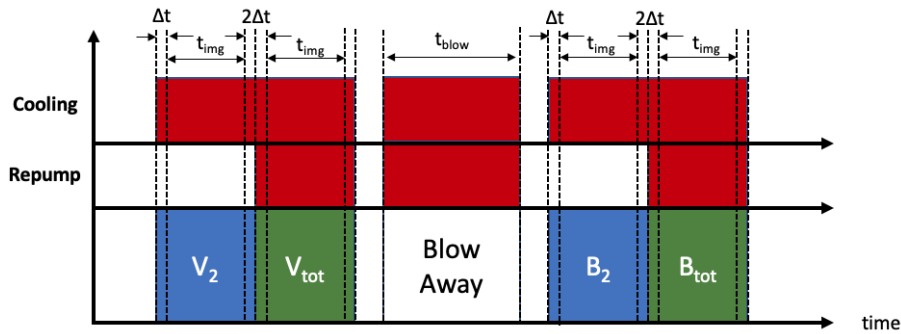


Figure 3.33: State-sensitive detection sequence

⁴The power and the detuning of the detection beams are optimised to have the maximal scattering rate from the atom. At the same time, the expansion of the cloud during detection is kept minimal to avoid loss of signal (see Ref. [62]).

probability P_2 of an atom being detected on $F = 2$:

$$P_2 = \frac{N_2}{N_{\text{tot}}} = \frac{V_2 - B_2}{V_{\text{tot}} - B_{\text{tot}}}. \quad (3.13)$$

It is worth noting that the photodiode signals are all averaged over a duration of $t_{\text{img}} = 300 \mu\text{s}$ to filter out high-frequency noises, with the signals during the first and the last $\Delta t = 100 \mu\text{s}$ omitted to avoid any exponential transient behaviour of the amplifier in the photodiode.

Liquid Crystal detection

We use only the two vertical MOT beams (Z_+ and Z_- beams) for these fluorescence detection steps, so we can ensure that the atom sees only the σ^+ polarisation required for the cycling transition. These beams are polarised σ^+/σ^- during the MOT loading stage, but are converted to σ^+/σ^+ for detection by a liquid crystal modulator in the path of the Z_- beam. This can switch rapidly between the polarisations for the MOT and the detection.

Figure 3.34 shows the experimental implementation of the modulator, as an upgrade for Z_- -MOT beam. The diverging beam coming out of the fibre goes through a linear

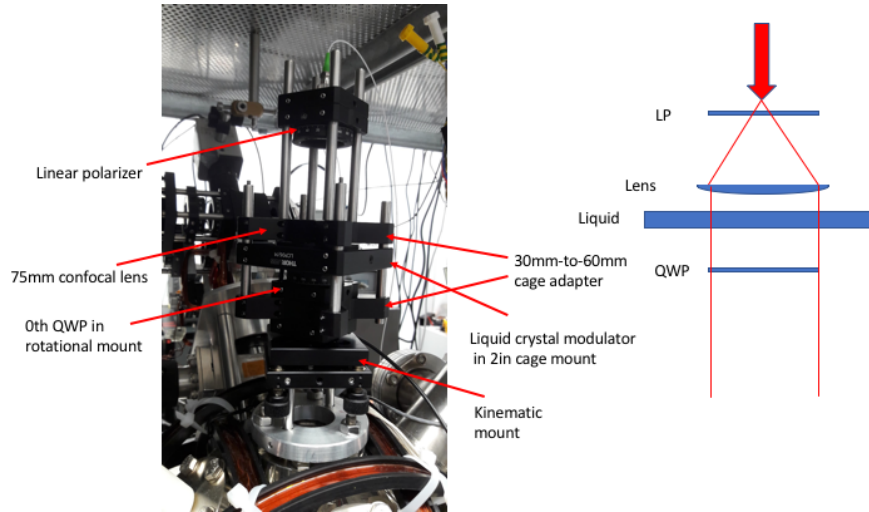


Figure 3.34: Upgraded Z_- -MOT beam with liquid crystal modulator: a picture of collimator with descriptions of components (left) and its schematic (right).

polariser (LP), which purifies the polarisation. A 7.5 mm lens collimates the light into the liquid crystal plate. The retardation of the crystal is controlled by the amplitude of a sinusoidal control voltage. We allow the modulator only to take two amplitudes, one, which does not affect the linear input polarisation, and the other, which rotates the polarisation by 90° . The switching between these two linear polarisations is triggered

by an external TTL signal. Finally, a zero-order quarter wave plate is aligned at 45° to generate the circular polarisations.

It is important to point out that the liquid crystal modulator switches the polarisation much more slowly in one direction than the other ($166\ \mu\text{s}$ versus $30\ \text{ms}$). The fast transition is used to go quickly from the detection of one atom cloud to the collection of the next, so that no additional time needs to be added to the sequence. The slow transition is accomplished at leisure during the interferometer measurement, while the MOT light is off.

3.8.2 Steps 1 and 2: Optical pumping

Step 1:

At the end of Sisyphus cooling, over 90% of atoms are left in the $F = 2$ state, while the rest remains in $F = 1$. Next, we use the MOT beams along the Z -axis to drive the $|F = 2\rangle \leftrightarrow |F' = 2\rangle$ transition with the repump light being switched-off. This pulse transfers all atoms from $F = 2$ to $F = 1$ in the $5^2S_{1/2}$ state via spontaneous decay. A bias field along the excitation beams is necessary to destabilise the dark states [67]. Figure 3.35 (a) shows the measured $F = 2$ populations (dots) as a function of the pulse

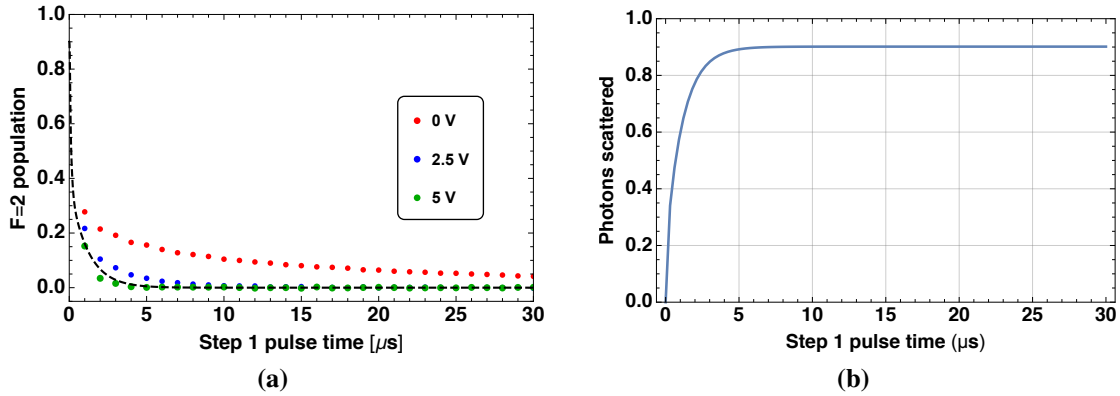


Figure 3.35: (a) exhibits the data of population on $F = 2$ after the first state preparation pulse. A faster population transfer is observed at a larger magnetic field along the Z -axis (given here in the units of control voltage of the current drivers). (b) simulates the number of scattered photons during the step 1 pulse as a function of the pulse duration. Around 0.9 photons are expected to be scattered by an atom, which corresponds to a temperature increase of $0.3\ \mu\text{K}$.

duration in step 1 for different bias field strength, which are plotted here in control voltage of the current drivers for the bias coils. After $5\ \mu\text{s}$ at $5\ \text{V}$ (corresponding to a bias field of roughly $1\ \text{G}$), we see that $F = 2$ population, which starts initially at approximately 0.9, depletes almost completely, because spontaneous emission has transferred almost all the population into the $(F = 1, m_F)$ states, which are dark. The time taken to pump out

$F = 2$ depends on the magnetic field along Z -axis, because this field destabilises the dark states in $F = 2$. In addition, we have simulated the dynamics of the populations of all magnetic sub-levels of $F = 2$ and $F = 1$ states during step 1. Using the Master equation⁵ with the known beam intensity of 3.7 mW/cm^2 , we vary the bias field in the Master equation to match the data collected at 5 V. The dashed black line in Fig. 3.35 (a) simulates the population on $F = 2$ with a bias field of 0.7 G, which is in good agreement with the data. By summing up all scattering rates from each Zeeman sub-level, we estimate the number of photons scattered during step 1 in Fig. 3.35 (b). The simulation shows that the number of scattered photons stagnates with vanishing population on $F = 2$. Close to 0.9 photon is scattered during step 1, which corresponds to an increase of temperature of around $0.3 \text{ } \mu\text{K}$. This increase exceeds the sensitivity of the time-of-flight measurement using the CCD camera.

Step 2:

After a $5 \text{ } \mu\text{s}$ step 1 pulse, more than 98% of the population occupy the $F = 1$ state. However, atoms are distributed, approximately evenly, among the three magnetic sub-levels, $m_F = 0, \pm 1$ of this manifold. Step 2 of the state preparation sequence transfers atoms from the $(F = 1, m_F = \pm 1)$ state into $(F = 1, m_F = 0)$ by driving the $|F = 1\rangle \leftrightarrow |F' = 0\rangle$ transition with a pair of counter-propagating σ^+/σ^- beams along the Z -axis of the experiment. This frequency is generated as a sideband from the laser beam used in the step 1, which drives $|F = 2\rangle \leftrightarrow |F' = 2\rangle$ transition, as mentioned in section 3.2.1 and section 3.7. Due to the setup of the MOT laser, μQuans , the carrier and sideband frequencies can't be separated in this step.

To measure the increase of population in $(F = 1, m_F = 0)$, we transfer all population from $(F = 1, m_F = 0)$ to $(F = 2, m_F = 0)$ with a resonant microwave pulse with almost 100% efficiency, so that they can be detected by fluorescence by driving the cycling transition, as described in section 3.8.1. (Section 3.8.3 describes the microwave source in more detail.) Figure 3.36 shows the population in $(F = 1, m_F = 0)$ as a function of the step 2 pulse duration. We find that 94% of atoms are concentrated in $(F = 1, m_F = 0)$ after optimising the frequency and the power of the sideband light. The remaining 6% are distributed between $(F = 1, m_F = \pm 1)$ states. This is due to the impure $\sigma^\pm$ polarisations of the $|F = 1\rangle \leftrightarrow |F' = 0\rangle$ frequency in the $Z_\pm$ beams. Any presence of π -polarisation is able to drive $(F = 1, m_F = 0)$ atoms back to $(F = 1, m_F = \pm 1)$ states. This is also incorporated into the Master equation to model the dynamics on $(F = 1, m_F = \pm 1)$ for impure step 2 pulse. In the model, we again adapt the experimentally obtained peak intensity of the $Z_\pm$ -beams of 3.7 mW/cm^2 and the same bias field of 0.7 G, obtained from

⁵The credit for the simulation goes to Malcolm Boshier from LANL.

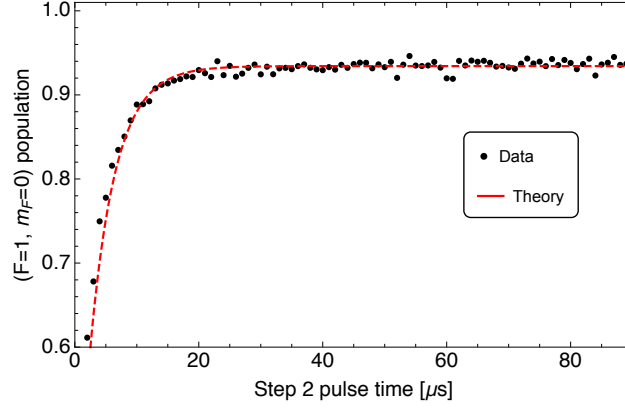


Figure 3.36: The measured ($F = 1, m_F = 0$) population is shown as a function of step 2 pulse duration in black dots, which is also simulated in red dashed line, by assuming that 5% of the total optical power is distributed in $|F = 1\rangle \leftrightarrow |F' = 0\rangle$ frequency and 1.5% of this is in the π -polarisation, which prevents complete population transfer.

the simulation in step 1. By varying the power distribution between the carrier frequency ($|F = 2\rangle \leftrightarrow |F' = 2\rangle$), its sidebands ($|F = 1\rangle \leftrightarrow |F' = 0\rangle$) in $\sigma^\pm$ and π polarisations, we understand the observed, imperfect state preparation in ($F = 1, m_F = 0$). The data in Fig. 3.36 are very well explained, as shown in red dashed line, by assuming that 5% of the total optical power is in the sideband, with 2.5% of that in the unwanted π -polarisation. Furthermore, Fig. 3.37 (a) shows the expected population in the ($F = 1, m_F = \pm 1$) states after a 60 μs step 2 pulse as a function of the intensity ratio between π and $\sigma^\pm$ -light ($I_\pi/I_{\sigma^\pm}$). The graph is plotted for three different power levels. From this, we conclude that the residual population in ($F = 1, m_F = \pm 1$) is more affected by imperfect polarisation than by inadequate power on the pumping transition.

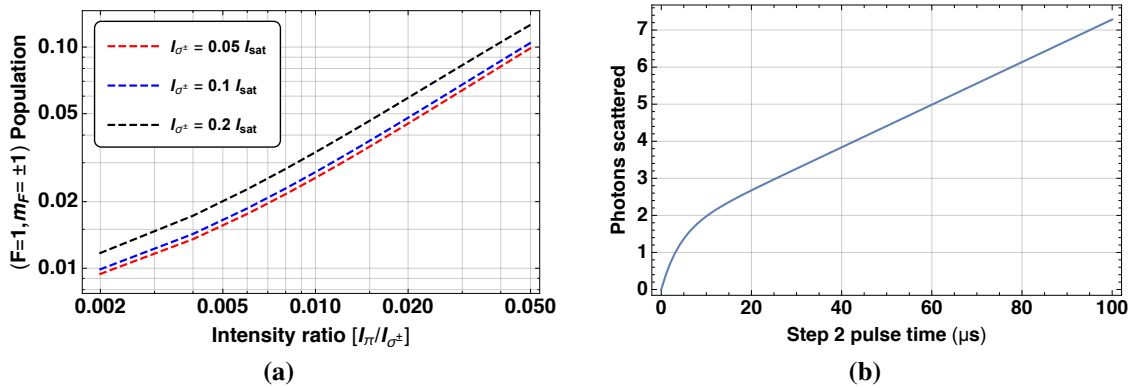


Figure 3.37: (a) illustrates the total residual population in the ($F = 1, m_F = \pm 1$) states as a function of the ratio between π and $\sigma^\pm$ -light intensities, $I_\pi/I_{\sigma^\pm}$. An intensity ratio of 2.5 % is sufficient to leave approximately 6 % of atoms remaining in $m_F = \pm 1$ manifolds after step 2. (c) shows the number of scattered photons during a step 2 pulse for $I_\pi/I_{\sigma^\pm} = 0.025$.

During the step 2 pulse, we expect additional heating of the cloud due to the presence of the more intense carrier frequency. Figure. 3.36 (b) exhibits the number of scattered

photons as a function of the step 2 pulse duration. For a $60\text{ }\mu\text{s}$ pulse, 5 photons are expected to be scattered from each atom, which correspond to an increase of temperature by $2\text{ }\mu\text{K}$. This has been measured in our device and we have observed an increase of temperatures on both axes from 7 to 8 and 6 to $9\text{ }\mu\text{K}$, respectively.

3.8.3 Step 3: Microwave transition

In step 3, we use a microwave pulse to transfer all the $(F = 1, m_F = 0)$ atoms to $(F = 2, m_F = 0)$ in preparation for the velocity-selective Raman pulse of step 4. This is a magnetic dipole transition, with the selection rules of $\Delta m_F = 0, \pm 1$. We apply a static magnetic field in order to pick out the desired transition using the Zeeman shift, and we optimise the transition rate by rotating the static field until it is parallel to the microwave magnetic field and perpendicular to microwave's axis of propagation. In switching the static field direction from the Z -axis to this new direction, the rotation should be slow enough compared with the Larmor frequency that no atoms are lost from $(F = 1, m_F = 0)$. From this perspective, we should require only a few microseconds. However, the eddy currents induced in the chamber wall die away much more slowly than that, therefore, a settling time of 4 ms is added to the sequence.

Setup

The C-band microwave horn (*Pasternack* (PE9859/SF-10)) is mounted outside the vacuum chamber with its exit located some 12 cm from the centre of the chamber, and with the magnetic field coarsely aligned to the Raman axis. Figure 3.38 illustrates the rf components. Two dual-channel microwave sources (*Windfreak*) are locked onto the 100 MHz clock reference from the μQuans laser (amplified twice and then split by a power splitter). These generate three rf frequencies that drive the (F, m_F) transitions: $(1, -1) \rightarrow (2, -1)$, $(1, 0) \rightarrow (2, 0)$ and $(1, +1) \rightarrow (2, +1)$ at frequencies f_{-1} , f_0 and f_{+1} respectively. These transitions are resolved by the application of a 1.4 G magnetic field. One *Windfreak* source puts out the single frequency f_0 , as required for step 3. The other provides us with $f_{\pm 1}$, which we use to clean up any atoms that have been left in $(F = 1, m_F = \pm 1)$ after step 2. TTL switches allow us to select either f_0 or $f_{\pm 1}$. The three outputs are combined using two power splitters. Because it is important not to expose the atoms to any microwave power while the interferometer is measuring acceleration, a master rf switch before the main amplifier kills off any residual power leaking through the primary TTL switches.

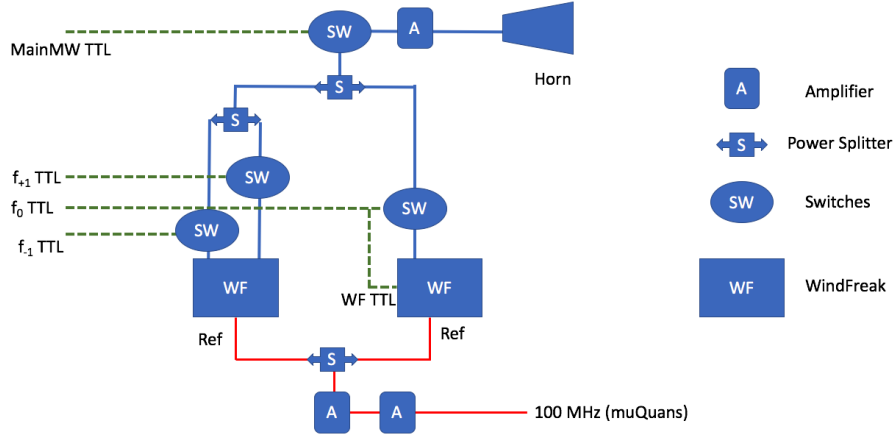


Figure 3.38: The microwave setup for step 3.

Microwave spectrum scan

The efficacy of step 2 is readily checked using this microwave setup. Figure 3.39 (a) shows the spectrum of (Zeeman resolved) microwave transitions ($m_F \rightarrow m'_F$) with the level scheme illustrated in (b), where the prime here indicates the $F = 2$ ground state. The signal is detected by the laser-induced fluorescence from $F = 2$ atoms. The red spectrum, taken in a bias field of 1.4 G after step 1, exhibits six resonances with $\Delta m_F = 0$ and $\Delta m_F = \pm 1$. Transitions with $\Delta m_F = 0$ are indicated by $(-1 \rightarrow -1')$, $(0 \rightarrow 0')$ and $(1 \rightarrow 1')$, whereas transitions with $\Delta m_F = \pm 1$ by $(-1 \rightarrow 0')$, $(0 \rightarrow -1')$, $(0 \rightarrow 1')$, $(1 \rightarrow 0')$ and $(1 \rightarrow 2')$. It is clear that $(\pm 1 \rightarrow \pm 1')$ and $(1 \rightarrow 2')$ peaks disappear after step 2 and $(\pm 1 \rightarrow 0')$ peak heights are reduced, as seen in the black spectrum. This shows that step 2 is effective in removing the atoms from $(F = 1, m_F = \pm 1)$. By contrast, the population at resonances excited from $(F = 1, m_F = 0)$ increased closely to 1. One does not see the 6% remaining on $(F = 1, m_F = \pm 1)$ here as a result of very sparse sampling. Figure 3.39 (c) shows the Rabi flopping as the duration of the microwave pulse is varied, with the frequency set on resonance with the $|0\rangle \rightarrow |0'\rangle$ transition. Although the oscillation decays, presumably due to the inhomogeneity over the atom cloud, the $F = 2$ population at the first π time is nevertheless around 94%, which corresponds to 100% of population initially in $(F = 1, m_F = 0)$.

In addition, by addressing the frequencies $f_{\pm 1}$ with microwaves, we could also drive atoms in $(F = 1, m_F = \pm 1)$ up to $F = 2$ to be cleared out by the blow-away beam, which could improve the visibility of the interferometer fringes that measure acceleration. However, we are not currently implementing that idea because the fluctuations of the ambient magnetic field change the atomic transition frequencies by more than the microwave line width. We are planning on characterising the fluctuation of the ambient magnetic field to design an appropriate compensation circuits, which can stabilise those

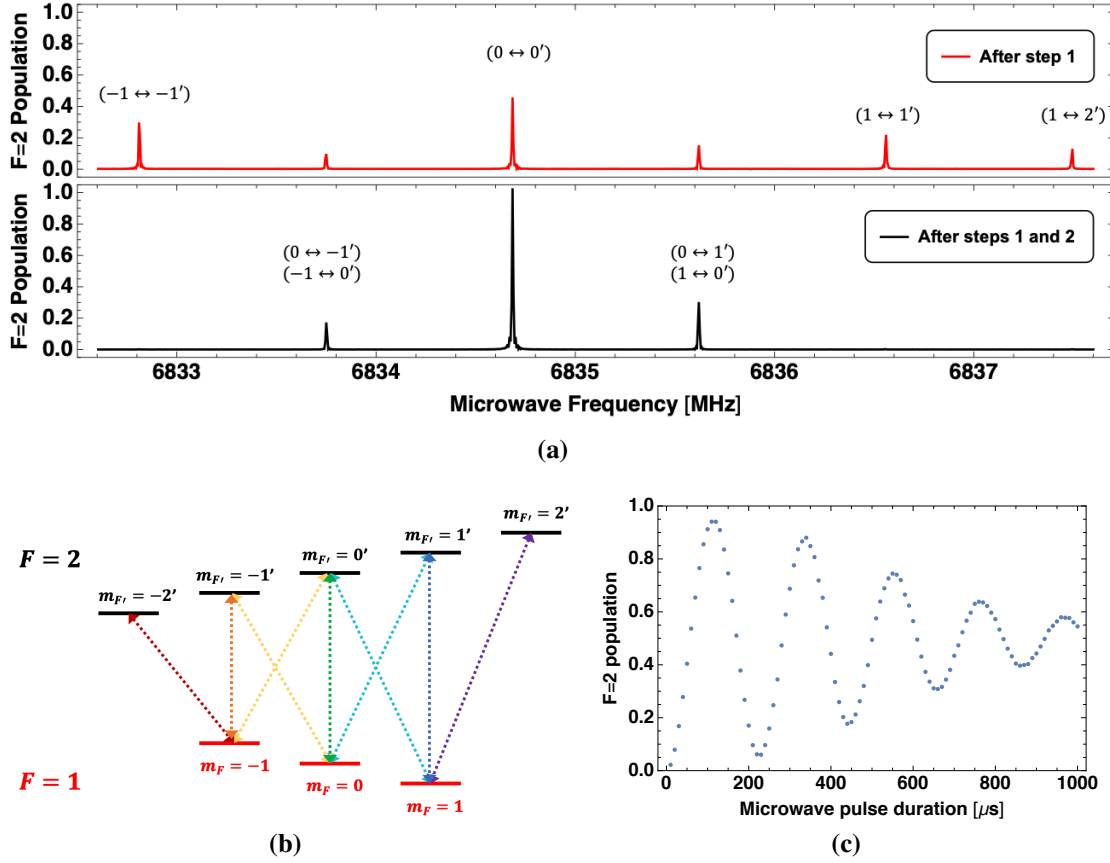


Figure 3.39: Microwave spectra (a). Red: Before step 2. Black: After step 2. The $(\pm 1 \rightarrow \pm 1')$ and $(1 \rightarrow 2')$ transitions vanish after step 2 pulse. (b) shows the Zeeman splitting of the 8 magnetic ground states, with transitions of the same frequency illustrated in the same colour. (c) is the microwave pulse duration scan on the resonance of the $(0 \rightarrow 0')$ transition.

fluctuations.

3.8.4 Step 4: Velocity selection

As the interferometer pulses are velocity-sensitive (see Chapter 5), only a sub-set of atoms within the $8 \mu\text{K}$ velocity distribution is addressed by the Raman pulse. To achieve atomic interference having a large fringe contrast, we preselect a narrow velocity class of atoms by driving the $|F = 2, m_F = 0\rangle \leftrightarrow |F = 1, m_F = 0\rangle$ transition with a long, velocity-sensitive Raman pulse (executed before the 3 interferometer pulses). The remaining atoms are then cleared out from $F = 2$ using a resonant light beam. Before describing this and other Raman pulses in detail, it is essential to describe the Raman optics, which I do next in Chapter 4.

Chapter 4

Raman Optics

The aim of the project is to build an interferometer to measure horizontal acceleration in the presence of normal vertical gravity and with an additional acceleration within ± 0.3 g along any directions. We aim to operate with an acceleration accuracy of 100 ng. To achieve this, we choose an interferometer time of $2T = 50$ ms. After 50 ms, atoms fall approximately 12 mm under gravity and move 4 mm at 0.3 g along any other directions. Taking the expansion of a 8 μ K cloud into consideration, atoms could situate anywhere within an area of approximately 1 cm² across the Raman beam. This puts constraints on the uniformity of the beam intensity and the laser phase within the area of interest.

As the atoms fall, the beam intensity at the centre of mass of the cloud varies. This could cause imperfect π - and $\pi/2$ -pulses in the interferometer, but could be compensated by adjusting the laser power or pulse duration accordingly. However, the intensity gradient across the atom cloud remains. Therefore, every atom experiences different laser intensity, causing inhomogeneity in the Raman Rabi frequency. To minimise the gradient, we require a Raman beam with a large spot size. Limited by the laser power available, we expand the Raman light into a collimated beam, which corresponds to a Gaussian beam with a spot size of 38 mm.

Additionally we wish to avoid an uneven wavefront, which could imprint different phases on different atoms in the cloud. To avoid unnecessary wavefront aberrations in our Raman beam, we use the state-of-the-art optics, with a surface flatness of $\leq \lambda/50$. Furthermore, all the optics are mounted inside the vacuum to eliminate phase noise from air turbulence.

In this chapter, I describe first the Raman laser system we purchased. Next, I describe the most recent Raman optics design and implementation in more detail.

4.1 Raman lasers - *M-Squared*

The Raman laser system is provided by *M-Squared lasers*. The two Raman frequencies are generated by two Ti:Sapphire lasers, a master and a slave, and are denoted by $f_{R,M}$ and $f_{R,S}$. Each laser has a linearly polarised output with a polarisation extinction of -20 dB or better. The two beams are orthogonally, linearly polarised and are overlapped before being coupled into a PM fibre¹, which delivers at its output approximately 1.5 W of optical power with a reduced PER of -15 dB. This fibre delivers the Raman light to the sensing head.

The master laser, pumped by a 15 W doubled YAG laser, is locked at -1.13 GHz red-detuned from the cooling transition $|F = 2\rangle \leftrightarrow |F' = 3\rangle$. This is the Raman detuning Δ shown in Fig. 2.2. A phase lock system is applied to lock the frequency of the slave laser $f_{R,S}$ to the master laser $f_{R,M}$ with an offset frequency, f_{Beat} , close to the atomic hyperfine transition $|1, 0\rangle \leftrightarrow |2, 0\rangle$ in the $5^2S_{1/2}$ fine structure level ($f_0 = 6.834682$ GHz). This controls the detuning δ shown in Fig. 2.2. To ensure that the transition probability resulting from each interferometer pulse is stable, the intensity fluctuations of the lasers need to be minimised and the phase of the microwave beat frequency needs to be well controlled.

Setting the intensity ratio

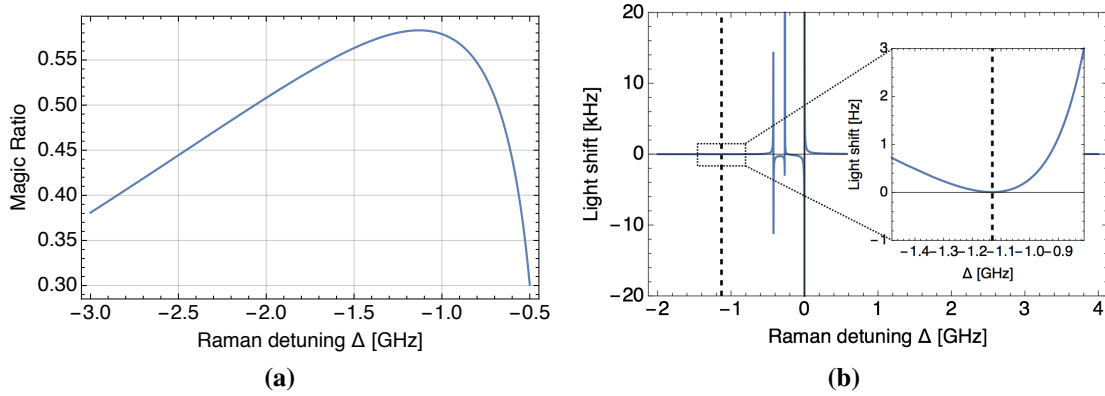


Figure 4.1: Calculation of the magic ratio versus the Raman detuning Δ (a). (b) exhibits the light shift of the clock transition as a function of Δ with an enlarged inset plot showing vanishing light shift around $\Delta = -1.13$ GHz.

Section 2.1 shows that the light shifts, i.e. AC stark shifts, are proportional to the laser intensities. In general, the shift of $|2, 0\rangle$ and $|1, 0\rangle$ differ, but with the detuning Δ

¹The input end of the Raman output fibre is accommodated with a Grin-collimator, which expands the incoming beam to protect the fibre tip from high intensity damages. This collimator is typically accompanied with a loss in PER.

that we have chosen, there is the possibility of making them the same by setting the right ratio of master and slave intensities. In that case the frequency of the hyperfine transition becomes insensitive to small variations of the laser frequency. Figure 4.1 (a) shows a calculation of this magic ratio as a function of the Raman detuning Δ , done by Prof. Hinds. We choose $\Delta = -1.13$ GHz, because the 0.58 magic ratio here is insensitive to small variations in the Raman detuning. Figure 4.1 (b) shows the light shift of the clock transition frequency as a function of the Raman detuning Δ with the intensity ratio set to 0.58. The light shift vanishes as expected for $\Delta = -1.13$ GHz, which is indicated by the black dashed vertical lines.

Phase lock stability

The master-slave laser phase lock has to be very good, as the phase difference between the counter-propagating light beams is directly picked up by the atoms at each interferometer pulse. Data² in the top section of Fig. 4.2 (a) shows the difference ϕ between the phase of the beat note and that of the reference oscillator, averaged over intervals of 0.2 ms. One can see a large drift spanning 30 radians. As the interferometer measures $\Delta\Phi = \phi_1 - 2\phi_2 + \phi_3$, this large shift is not in itself a problem for the interferometer. Using the data provided in the top section of the Fig. 4.2 (a), I approximate the interferometer phase $\Delta\Phi$ for interferometers operating at $T = 0.2, 5,$ and 25 ms in the bottom section of the Fig. 4.2 (a). Here, $\Delta\Phi(t)$ at time t is calculated using three data points, $\phi(t - T)$, $\phi(t)$ and $\phi(t + T)$, where every $\Delta\Phi$ is estimated at a rate of $1/2T$. From shot to shot, $\Delta\Phi$ is distributed around 0 rad with a standard deviation of approximately 0.1 rad. (The spread of $\Delta\Phi$ for $2T = 0.4$ ms in the black line looks broader than the red and green lines due to higher number of data points.) This estimated interferometer phase $\Delta\Phi$ over time is then converted into the interferometer phase noise $\sigma_{\Delta\Phi}$ as a function of numbers of measurements, as shown in the Allan deviation plot in Fig. 4.2 (b). This shows that $\sigma_{\Delta\Phi}$ noise will be increasingly suppressed if more measurements are performed. Furthermore, Fig. 4.2 (c) exhibits the interferometer phase noise as a function of the interferometer time $2T$ after 1 s of measurement. This demonstrates that the laser system provided by *M-Squared Lasers* is able to reach an interferometer phase noise of just ≤ 10 mrad, corresponding to an acceleration uncertainty of 100 ng after 20 measurements for an interferometer operating at $2T = 50$ ms (see section 7.1.3 for further discussion about noise). However, to observe the limitation of the stability of the Raman laser beat note phase, more data of the laser beat note phase are required.

It is worth noting that every point of ϕ in the top section of the Fig. 4.2 (a) is averaged over 0.2 ms, which is significantly longer than the durations of π - and $\pi/2$ -pulse (typi-

²The raw data are provided by Joseph Thom and Stephen Webster at *M-Squared Lasers*.

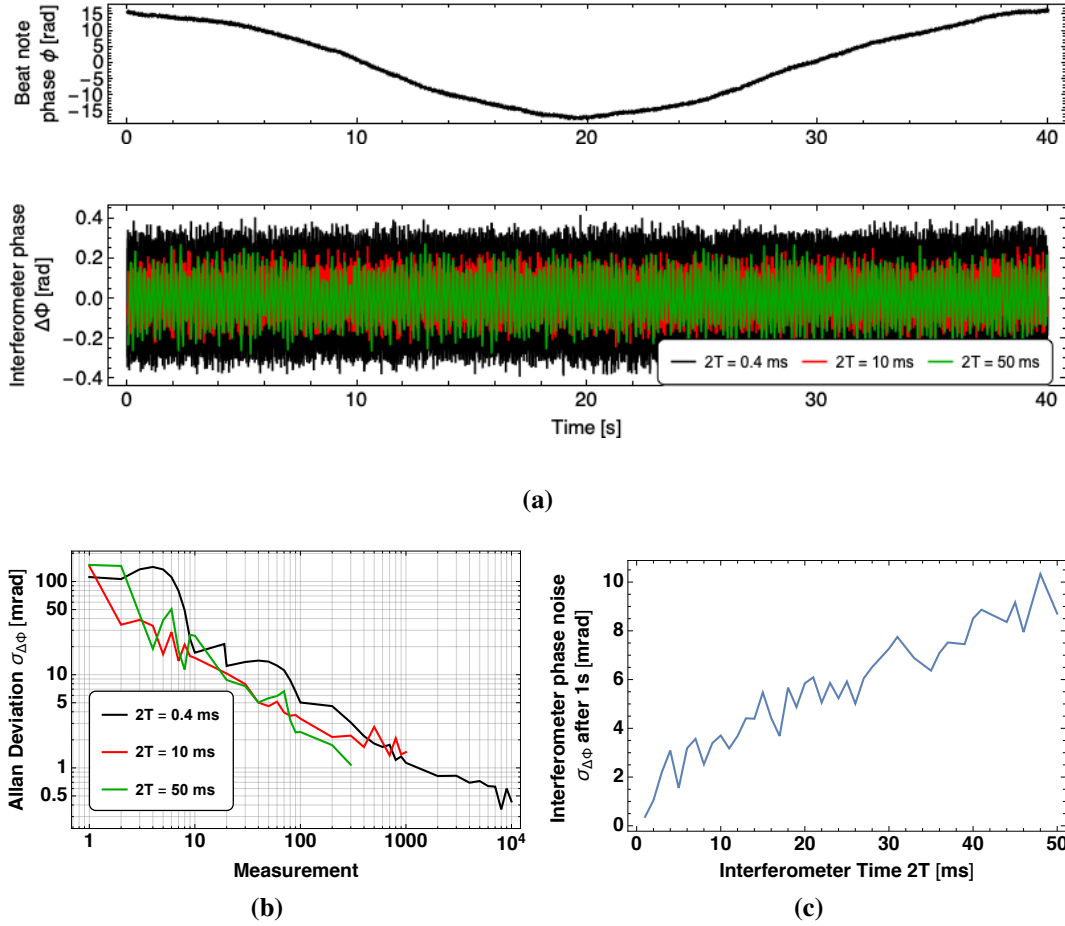


Figure 4.2: Characterisation of the *M-Squared Laser* system: (a) shows the difference ϕ between the Raman laser beat note phase and a reference oscillator (top). The interferometer phases $\Delta\Phi$ for $2T = 0.4, 10$ and 50 ms are estimated using the data in the top section in black, red and green lines respectively. (b) shows the Allan deviation of the interferometer phase noise $\sigma_{\Delta\Phi}$ as a function of the number of measurements. (c) plots the estimated interferometer phase noise after 1 s measurement for various $2T$ interferometers.

cally $\leq 15\mu\text{s}$). Hence, the bandwidth of the approximated interferometer pulses is also smaller than what the atoms would experience. Nevertheless, it provides us with a good first estimate of the stability of the interferometer phase due to the phase drift of the laser beat note.

4.2 In-vacuum Raman collimator

The optical setup consists of two separate parts, as shown in Fig. 4.3, a collimator and a mirror assembly, which also includes a state-of-the-art high bandwidth pendulous accelerometer to record vibrational noise (see Chapter 7). This section explains the design and the setup of the collimator.

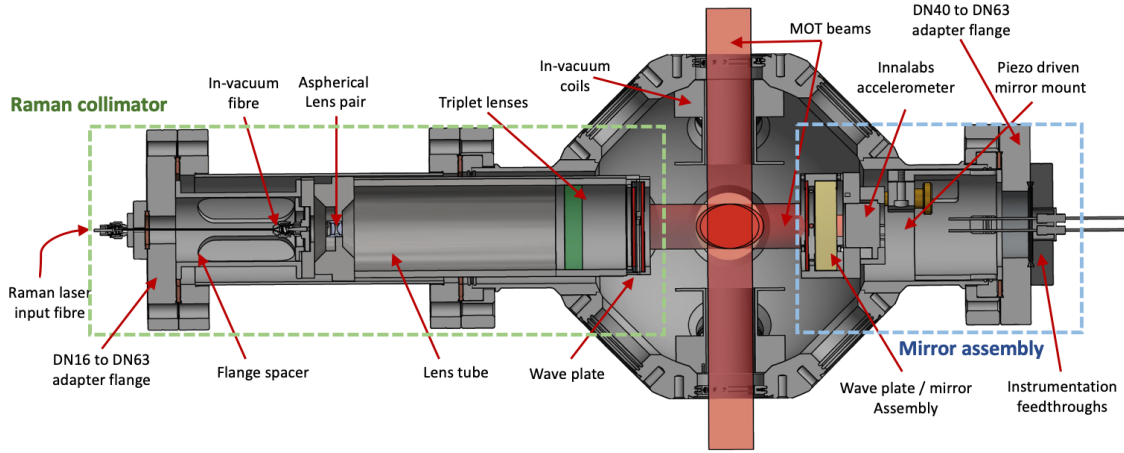


Figure 4.3: Schematic of the Raman optics design. Collimator (surrounded by green lines): An adapter piece houses the pair of aspheric lenses, which connects the fibre plate with the lens tube. A rotation mount for the waveplate is mounted on the other side of the lens tube together with a set of triplet lenses. Mirror assembly (surrounded by blue lines): This unit allows us to firmly fix the pendulous accelerometer behind the mirror.

4.2.1 Design

A single-mode polarisation-maintaining fibre guides two orthogonally, linearly polarised frequencies from the *M-Squared* laser system into the Raman collimator. The optics expands and collimates the light to form the large beam described in section 4.2.2. A high-quality quarter-wave plate prepares the linear polarisations into a left- and a right-handed circularly polarised light, i.e. driving $\sigma^\pm$ -transitions in an atom in the presence of a magnetic field along the laser propagation axis. An ideal collimated Raman beam with a uniform intensity distribution and a plane wave front assures that all atoms undergo the same transition and pick up the desired laser phase difference, while moving across the laser beam. However, in a real system, the intensity and hence the Rabi frequency varies across the atom cloud, which leads to inhomogeneity of the $\pi/2$ - and π -pulses. As a result, this decreases the contrast of the interference fringe and affects the sensitivity of the interferometer. We can reduce the loss of fringe contrast by keeping the atoms close to the centre of the Raman beam during the interferometry, where the beam intensity varies the least. For an interferometer operating at $2T = 50$ ms, vertically launching the atoms against the gravity at a velocity of 25 cm/s (see section 3.7.1) keeps the centre of mass of the atom cloud approximately 3 mm around the beam centre. To further illustrate this, I use Eq. 2.16 and Eq. 2.18 and relate the single-atom interference fringe contrast C as a function of the Rabi frequency $\Omega(r_i) \propto \exp(-r_i^2/w^2)$ during i -th interferometer pulse by

$$C(\Omega(r_1), \Omega(r_2), \Omega(r_3)), \quad (4.1)$$

where r_i is the radial position of the atom across a Raman beam with a waist of w . For simplicity, I average in one dimension across a Gaussian atom density with standard deviation σ , the averaged fringe contrast $\tilde{C}$ can be described by

$$\tilde{C} = \int n_0 e^{-\frac{\rho^2}{2\sigma^2}} C(\Omega(r_1 - \rho), \Omega(r_2 - \rho), \Omega(r_3 - \rho)) d\rho, \quad (4.2)$$

where n_0 is the peak atom number density. Figure 4.4 plots the Eq. 4.2 as a function of Raman beam spot size ω . The red line shows an atom cloud having $\sigma = 2.5$ mm, which falls across the beam under gravity within a $2T = 50$ ms interferometer. To keep the loss of the fringe contrast below 10%, a Raman beam with a spot size of > 30 mm is required. When atoms are launched vertically against the gravity, the fringe contrast significantly

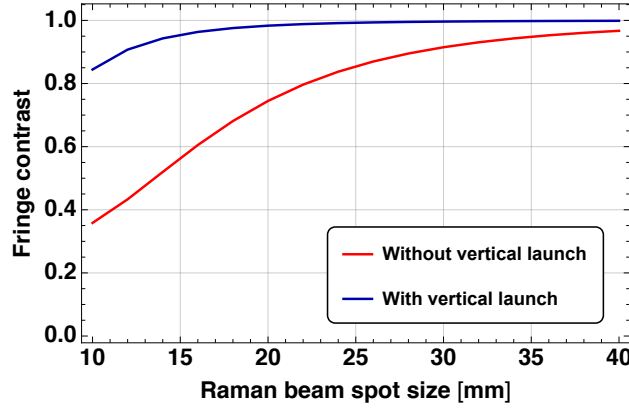


Figure 4.4: Estimated fringe contrast is shown as a function of Raman beam spot size. The calculation is performed on an atom cloud having $\sigma = 2.5$ mm falling across the beam under gravity perpendicular to the optical beam axis in red. The blue line exhibits the fringe contrast for atoms, which are additionally launched vertically against the gravity at a velocity of 25 cm/s.

improves, shown in the blue line. The fringe contrast remains almost constant for Raman beams with a spot size of > 20 mm.

The current Raman collimator is based on an existing collimator, which was initially designed with the help of Prof. Peter Torok at Imperial College London. The existing collimator benefits from the simplicity of expanding the light from the fibre tip using a single triple lens system. The triplet lens was customised by *IC Optical Systems Ltd.* to provide a surface roughness of better than $\lambda/25$ (at 780 nm). Table 4.1 summarises the prescription for individual lens surface of the high-quality triplet lens, which are illustrated in Fig. 4.5 (a).

4.2. IN-VACUUM RAMAN COLLIMATOR

Surface	Radius [mm]	Distance [mm]	Glass	Diameter [mm]
1	183.855	14.891	N-F2	50
2	-95.55	8.892	N-SF15	50
3	-941.25	1.799641	Air	
4	85.955	6.922	N-F2	50
5	210.265	123.3573	Air	

Table 4.1: Prescription for the triplet lens system provided by *IC Optical Systems Ltd.*: Each surface is characterised by the radius of curvature, the distance to the next surface, glass type and its diameter.

As this collimator outputs a beam with a spot size of 12 mm, we extended it with an off-the-shelf mounted aspheric lens pair (*C230260P-B*), purchased from *Thorlabs Inc.* The aspheric lens pair with focal lengths of 15.29 mm and 4.51 mm increases the divergence of the light incident on the triplet, thereby increasing the (calculated) spot size to 43 mm. Figure 4.5 (a) shows a schematic of the current collimator design. Figure 4.5 (b)

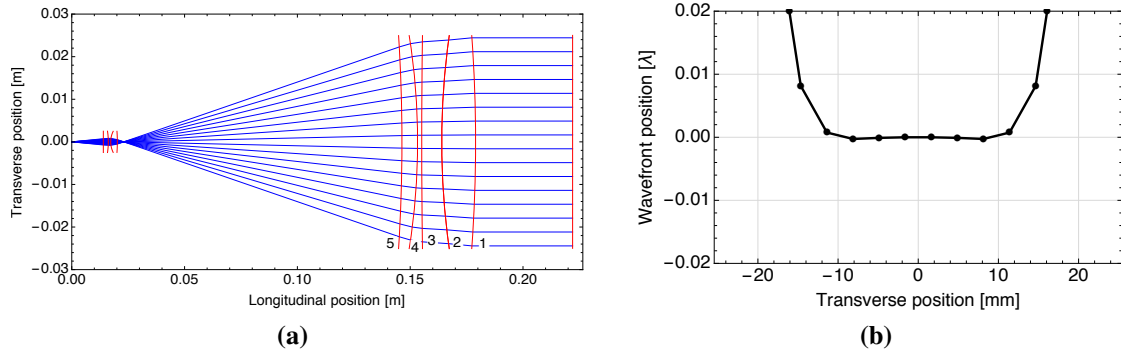


Figure 4.5: Design of the Raman collimator: (a) illustrates the aspheric lenses and the triplet lens system with ray tracing method. (b) the wavefront position is calculated using ABCD matrices transversely across the beam.

shows the position of the wave front in the units of wavelength, $\lambda = 780$ nm, as a function of transverse position across the beam. This was calculated using ABCD matrices. For perfect optics, the simulation shows that the current optical collimator should deliver a wave front flatness better than $\lambda/100$ over a circle almost 3 cm across.

Collimation requirement

The collimation of the Raman beams is important to ensure an accurate measurement of the acceleration, as we now discuss. The phase of each Raman beam can be described by assuming a Gaussian beam within the paraxial approximation

$$\phi_i^{\text{fr}}(r, z) = k z + k \frac{r^2}{2R(z)} - \arctan\left(\frac{z}{z_R}\right), \quad (4.3)$$

where subscripts “f/r” indicate the forward and the retro-reflected beams respectively, z_R is the Rayleigh length and $R(z)$ is the radius of curvature as a function of longitudinal position z . Under gravity, the positions of the falling atoms during the interferometer pulses are described by the increasing cylindrical radial coordinate r . For simplicity, we ignore kz and the Gouy phase $\arctan\left(\frac{z}{z_R}\right)$ because these do not depend on r . Therefore, the phase of the beat note is given by

$$\phi_i(r, z) = \phi_i^f(r, z) - \phi_i^r(r, z + 2L) \approx k \frac{r^2}{2} \left(\frac{1}{R(z)} - \frac{1}{R(z + 2L)} \right), \quad (4.4)$$

where L is the distance between the atom and the mirror. By inserting Eq. 4.4 into Eq. 2.23, the interferometer phase due to variation of radii of curvature can be expressed as a function of divergence angle, as a measure of beam collimation. Figure 4.6 shows the interferometer phase for atoms falling under gravity in a $2T = 50$ ms interferometer. Here, the divergence angle is also related to a more experimentally accessible quantity,

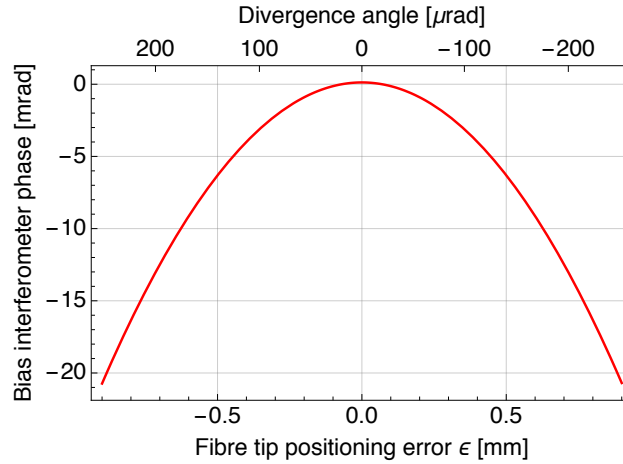


Figure 4.6: In a static environment, an imperfectly collimated beam induces a bias in the interferometer phase, which is shown here as functions of the Raman beam’s divergence angle and of the positioning error ϵ of the fibre tip. The calculation is performed for atoms falling under gravity during a $2T = 50$ ms interferometer.

the position error ϵ between the fibre tip and the aspherical lens pairs, whereas the separation between the aspherical lens pairs and the triplet lens system is assumed to be ideal. Under these assumptions, it shows that the Raman beam needs to be collimated with a divergence angle between $\pm 175 \mu\text{rad}$ to produce an additional bias interferometer phase of ≤ 10 mrad (corresponds to $1 \mu\text{m/s}^2$ accuracy in the acceleration measurement for an interferometer operating at $2T = 50$ ms).

In a static environment, where the trajectory of the atoms under gravity is the same in each measurement cycle, the interferometer phase contribution due to the radii of curvature is constant. However, in a dynamic environment with continuously varying un-

known acceleration between $\pm 3 \text{ m/s}^2$, the atom cloud can fall into any position in the beam within a diameter of around 15 mm in a 50 ms interferometer. The interferometer effectively also probes the wavefront of the Raman beam within the area covered by the atom cloud, apart from the acceleration. Hence, we treat the contribution due to radii of curvature as a source of varying systematic error in the acceleration measurement. For this reason, it is desirable to have a well collimated Raman beam to minimise the change of the radii of curvature.

4.2.2 Setup

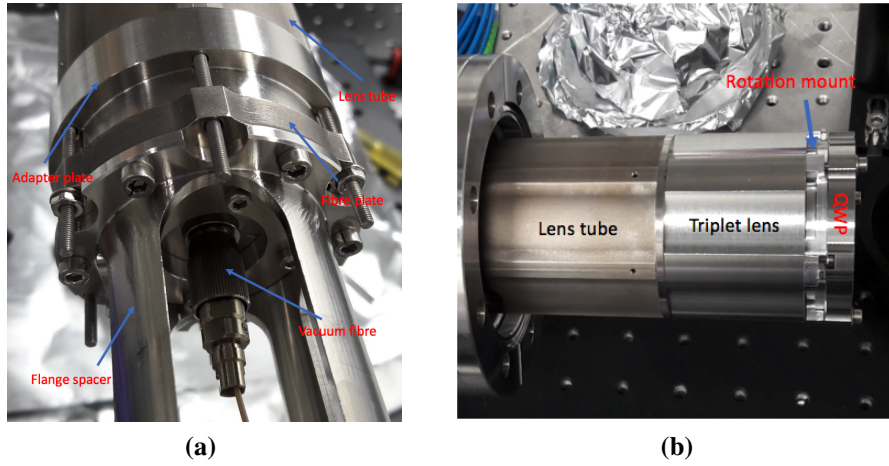


Figure 4.7: Raman collimator assembly inventory schematic (a). (b) and (c) show two photographs of parts of the collimator during assembling.

The section surrounded by the green lines in Fig. 4.3 shows a schematic of the Raman collimator assembly. The assembly is mounted on a DN16CF-to-DN63CF adapter flange, which is connected to the main chamber via an extension DN63CF full nipple with a length of 16 cm. A vacuum fibre feedthrough on a DN16CF flange is purchased from *Diamond SA*, in order to guide the light into the vacuum. The feedthrough is accommodated with a FC/APC female fibre connector on the air side and a 30 cm long, FC/APC male terminated fibre connector on the vacuum side allowing flexibility in the alignment. The bare fibre itself is coated with Kapton for low outgasing. The vacuum fibre feedthrough is attached to the DN16CF socket of the adapter flange. A flange spacer with a length of 71 mm connects the DN63CF adapter flange with a fibre adapter plate, onto which the male end of the vacuum fibre is attached, shown in Fig. 4.7 (b). A second adapter piece with a thickness of 22.68 mm is customised to house the aspherical lens pair. Both adapter plates are then connected to a 114.5 mm long lens tube, which overlap the focus of the aspherical lens pairs and the triplet lens, as shown in Fig. 4.7 (b).

At the output of the triplet lens, I have designed a rotation mount, which, on one hand, locks the triplet lens to the lens tube and, on the other hand, it allows adjustment to the orientation of a 2-inch high-quality quarter wave plate, as shown in Fig. 4.7 (c). The wave plate is purchased from *LightMachinery Inc.* with a specified phase retardation accuracy of $< \lambda/1000$. All components are manufactured for 2-inch optics with stainless steel for low outgasing.

Collimation alignment

Section 4.2.1 has described that a divergence angle within $\pm 175 \mu\text{rad}$ needs to be achieved to keep the acceleration noise below $1 \mu\text{m/s}^2$ in a $2T = 50 \text{ ms}$ interferometer. The collimation of the beam is more sensitive to the distance between the two sets of lenses than the distance between the fibre tip and the aspherical lens pair. Hence, we rely on the precision of the machinery for the lens tubes to set the distance between the lenses. The fine adjustment of the collimation is carried out by adding multiple shim plates (with a thickness within $200 - 300 \mu\text{m}$ each) between the fibre and the aspherical lenses. For each added shim plate, we project the beam onto a screen and record the spot size of the projected beam profile at regular intervals to determine the divergence of the beam.

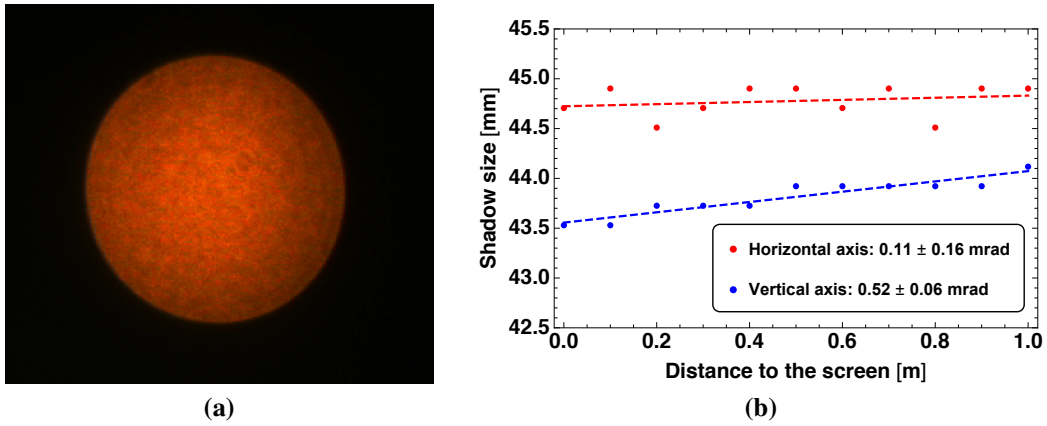


Figure 4.8: Image of the Raman beam with an effective waist of 38 mm on a screen is shown in (a). (b) summarises the divergence of the beam by looking at the shadow size of the clipped beam.

Figure 4.8 (a) is a picture of the collimated Raman beam projected on a screen. The rough intensity structure is predominately due to the surface quality of the aspherical lens pair, which is formed from molten glass. The speckles due to the roughness of the screen has a significantly higher spatial frequency than that of the intensity structure in (a). To estimate the spot size of the beam, I sum over the beam intensity over every 5 rows/columns³ of pixels in the picture, which are fitted by a parabola to extract the

³The intensity summed over 5 rows/columns are just sufficient for further data analysis. Simultaneously, this selection is small enough not to change the intensity profile.

horizontal/vertical spot sizes. From the full width of half maximum w_{FWHM} obtained from the parabola fits, I assign a Gaussian $1/e^2$ spot size w_0 for each sets of pixel rows/columns using

$$w_0 = \frac{w_{\text{FWHM}}}{\sqrt{2 \ln(2)}}. \quad (4.5)$$

An overall Gaussian spot size over the entire image is estimated by the mean of w_0 obtained from every 5 rows/columns of pixels, which gives a horizontal and vertical spot size of 40 and 35 mm respectively, indicating slight ellipticity in the beam. Due to the large error bar of few mm in the spot size estimation, I use the shadow of the clipped beam to estimate the divergence angle, which is more accurate due to a higher signal-to-noise ratio. Figure 4.8 (b) shows the variation of the shadow size as a function of the distance between the collimator and the screen. This method yields a sufficiently small divergence of 0.11 ± 0.16 mrad on the horizontal axis, fulfilling the design goal. However, the beam diverges with 0.52 ± 0.06 mrad on the vertical axis⁴. Comparing with Fig. 4.6, this corresponds to an interferometer phase of 90 mrad. This again can be treated as a varying source of systematic error in the acceleration measurement in environments with varying unknown acceleration.

4.3 Retro-reflecting mirror assembly

4.3.1 Setup

The second part of the Raman optics is the retro-reflecting mirror assembly. Figure 4.9 shows a photograph (left) and a schematic CAD drawing (right) of our Raman mirror assembly.

A hollow cylindrical tube (mirror holder) connects the ultra-high vacuum compatible kinematic mirror mount, purchased from *Radiant Dyes Laser & Accessories GmbH*, to a DN63CF-to-DN40CF adapter flange. The mirror mount is equipped with a pair of piezo transducers stacked together with mechanical actuators, which allow us to achieve high precision in the retro-reflection of the beams. The mechanical actuators are used for coarse adjustment in air, whereas the piezo transducers (with a maximal tilting angle of 1.2 mrad) are used for in-vacuum fine-tuning. An adapter plate, which is secured on the mirror mount, also connects the mirror/waveplate assembly with a UHV-compatible Quartz pendulum accelerometer, to directly track the movement of the mirror. All components are carefully chosen and tested for UHV usage.

The electric power and control wires are coated with Kapton, and the current from

⁴It is unclear where the astigmatism comes from. Further study is required to understand this.

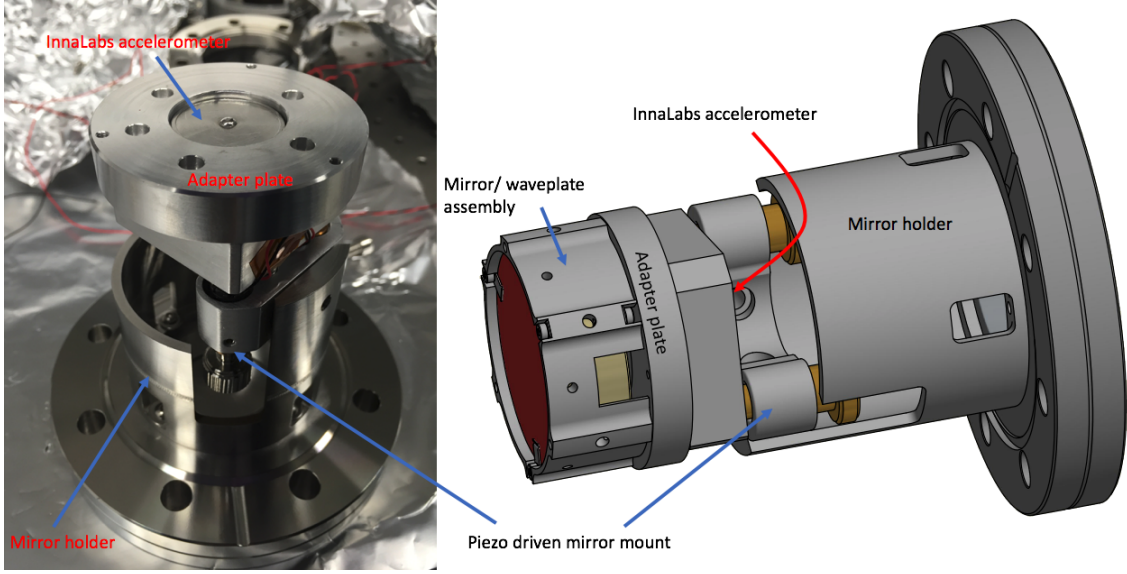


Figure 4.9: Raman mirror assembly: The picture on the left hand side exhibits the setup without the mirror/waveplate assembly.

the Quartz accelerometer is guided through a twisted pair of shielded Kapton wires, to minimise the pick-up of ambient magnetic and electric field fluctuations. A customised DN40CF flange with a C9 connector and two SMA feedthroughs enables us to read and control the accelerometer and mirror from outside the chamber.

Wave plate and mirror

In addition to the quarter wave plate at the output of the collimator, a second quarter wave plate is placed in front of the mirror to invert the polarisations of the retro-reflected beams with respect to the in-coming beams ($\sigma^+ \rightarrow \sigma^-$ and vice versa). As the retardation and surface uniformity of the quarter wave plate and mirror are directly probed by the atoms, we require the highest commercially available quality for these optics to minimise the systematic error in the interferometer. Figure 4.10 (a) shows a CAD inventory model of the assembly, customised by *LightMachinery Inc.* The quarter wave plate and a mirror are mounted together in a stainless steel holder with three optic retainers fixed by M2 screws on the side. A 2-inch mirror with 93% reflectivity at 780 nm is polished to a flatness of 29 nm over the entire surface. Figure 4.10 (b) is a contour plot of the relative mirror surface flatness in nm, which is shown as a function of mirror's horizontal and vertical positions (data provided by *LightMachinery Inc.*). The data are measured in a step size of 2 mm across the entire surface. We average the data over discs of different diameters around the centre of the mirror to estimate the standard deviation of the flatness. This is shown in Fig. 4.10 (c). In environments, where the acceleration can vary within ± 0.3 g in any direction additionally to the gravity, we expect the atoms to be within ± 8 mm of

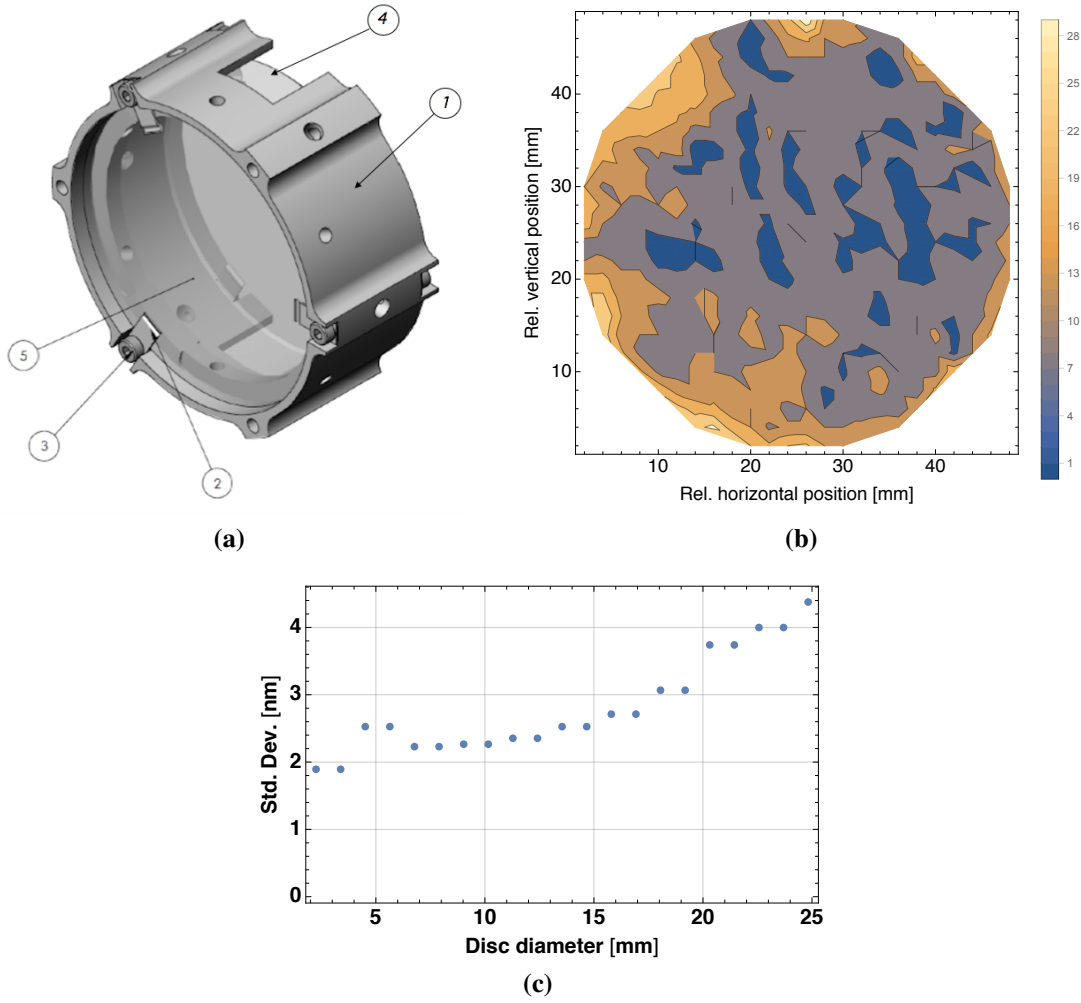


Figure 4.10: Waveplate/mirror assembly CAD drawing (a): 1. Stainless steel holder, 2. Optical retainer, 3. M2 screws, 4. Mirror glass and 5. Quarter waveplate. (b) surface profile mapping of the retro-reflecting mirror in nm (data provided by *LightMachinery Inc.*). (c) Standard deviation of the mirror surface flatness as a function of the diameter of the area disc averaged over.

the optical axis of the mirror. Over that region, the uncertainty in the position of the mirror surface has a standard deviation of 2.8 nm according to Fig. 4.10 (c). In addition, the thickness uniformity of the wave plate is measured to be 19 nm from peak to valley, which corresponds to an optical path length of 4.8 nm. These uncertainties add up in quadrature to be 5.6 nm and affect only the path length for the retro-reflected beam, k_2 . The laser beat note phase seen by the atoms at each interferometer pulse is described by $\phi_i = k_1 x + k_2(x + 2L_i)$, where L_i is the distance between the atom and the mirror at i -th pulse. Hence, the contribution to the interferometer phase due to the propagation over this path is

$$\delta\Phi_m = 2k_2(L_1 - 2L_2 + L_3). \quad (4.6)$$

Assuming that the uncertainties of L_i , σ_{L_i} are uncorrelated, I can estimate the systematic noise $\sigma_{\delta\Phi_m}$ of the interferometer phase due to the surface quality of the wave plate and the mirror using error propagation in Eq. 4.6. This yields a systematic noise of $\sigma_{\delta\Phi_m} = 221$ mrad within an area of interest with a diameter of 16 mm on the mirror.

However, when considering the impact of the mirror flatness on the whole atom cloud, we have to distinguish between the inhomogeneous variation of L over the cloud and the homogeneous variation, as the atoms move across the mirror. Due to the inhomogeneous variation of L , the interferometer phase varies across the cloud, which in turn results in a reduction of the overall fringe contrast. We approximate this by assuming $\delta\Phi_m$ to be randomly distributed around the mean value $\delta\bar{\Phi}_m$. The averaged probability of finding atoms in $F = 2$, $\tilde{P}_2$, after the interferometry can be described by

$$\begin{aligned}\tilde{P}_2 &= \int n_{\delta\Phi_m} \exp\left(-\frac{(\delta\Phi_m - \delta\bar{\Phi}_m)^2}{2\sigma_{\delta\Phi_m}^2}\right) \sin^2\left(\frac{\Delta\Phi + \delta\Phi_m}{2}\right) d\delta\Phi_m \\ &= \frac{1}{2} \left(1 - \exp(-\sigma_{\delta\Phi_m}^2/2) \cos(\Delta\Phi + \delta\bar{\Phi}_m)\right),\end{aligned}\quad (4.7)$$

where the $\sin^2$ -term is the single-atom interference fringe pattern related to Eq. 2.18. The other terms in the integral describe the normalised interferometer phase distribution due to the surface flatness with $n_{\delta\Phi_m}$ as the normalisation factor. This equation shows that the overall fringe contrast is $C = \exp(-\sigma_{\delta\Phi_m}^2/2)$ and for $\sigma_{\delta\Phi_m} = 221$ mrad, the reduction of the contrast to 0.98 is negligible. On the other hand, $\delta\bar{\Phi}_m$ causes a shift in the fringe pattern, which introduces an error in measuring the zero acceleration. In a dynamic environment, where the acceleration can vary within ± 0.3 g in any direction, $\delta\bar{\Phi}_m$ varies according to the trajectory on the mirror, which the atoms move along. This is interpreted as a systematic error in the accuracy of the measured acceleration. Figure 4.11 shows the calculated interferometer phase shift $\delta\bar{\Phi}_m$ as a function of the transverse acceleration (the horizontal axis in Fig. 4.10 (b)), which is perpendicular to the gravity (the vertical axis in Fig. 4.10 (b)) and the Raman beam axis. For the calculation, a cloud with a diameter of 2.5 mm is assumed to fall under gravity in a $2T = 50$ ms interferometer. Any random transverse acceleration within ± 0.3 g causes the interferometer phase to vary within a standard deviation of 24 mrad, which corresponds an acceleration noise of $2.4 \mu\text{m/s}^2$. This provide us with an estimate of the lower boundary for the sensitivity of the interferometer in a full dynamic environment.

Quartz pendulous accelerometer

A quartz-pendulum MEMS accelerometer measures the motion of the retro-reflecting mirror. This has sufficient accuracy to identify the order number of the atom interferom-

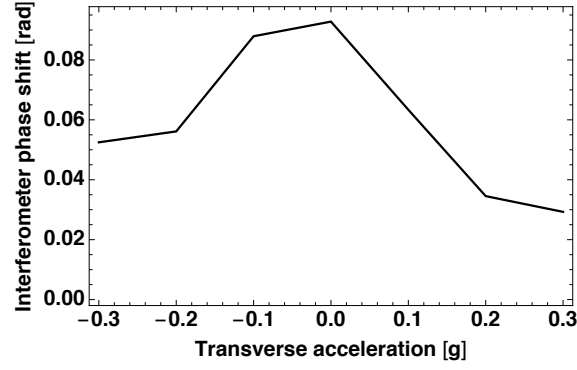


Figure 4.11: The interferometer phase shift is calculated for a 5 mm cloud falling under gravity using the data presented in Fig. 4.10 (b). A transverse acceleration between ± 0.3 g is introduced to observe the shift in the fringe pattern, which varies within a standard deviation of 24 mrad for a $2T = 50$ ms interferometer.

eter fringe and also provides acceleration information during the dead time of the atom interferometer, which is necessary for continuous operation in real world navigation. Figure 4.12 (a) illustrates the accelerometer, together with its working principle shown in (b). The proof mass is a quartz pendulum (“flapper”) with a pair of coils situated be-

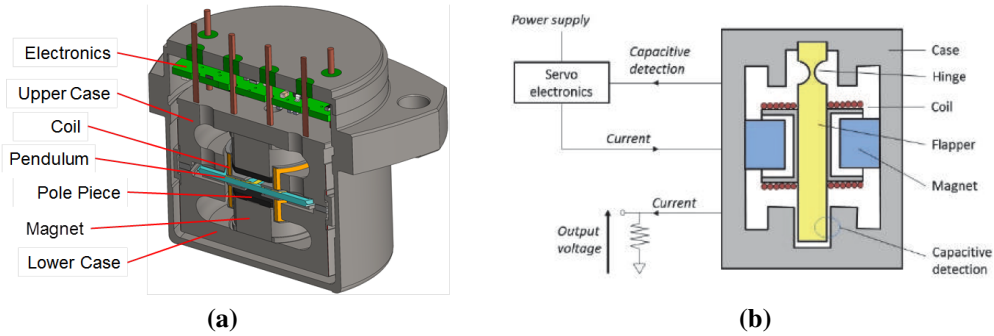


Figure 4.12: A schematic of a Quartz pendulous accelerometer is shown in (a), whereas (b) illustrates its working principle. Figures taken from Ref. [5]

tween a pair of magnets. When acceleration is applied, servo electronics derive an error signal from the capacitive detection, and this is used to generate a current through the coils in order to keep the proof mass in its null position with a frequency bandwidth of up to 2 kHz. The same restoring current is used in combination with an external resistor to generate an output voltage. The *Innalabs* (AI-Q-2010) that we use has a scale factor specified to be 1.235875 mA/g with a temperature variation of -29.46 ppm/ $^{\circ}\text{C}$ and a bias of -2.1 mg, varying by 8.2 $\mu\text{g}/^{\circ}\text{C}$.

It is important for us to know the voltage-converting resistance to a high accuracy. We have measured the value $5981.1 \pm 2.3 \Omega$ using a calibrated current source, as shown in Fig. 4.13. This value will be referred to in the remaining thesis. A more thorough description of the integration of the quartz accelerometer into atom interferometer will be covered in

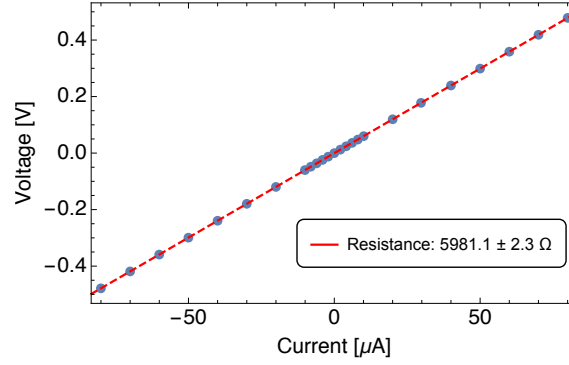


Figure 4.13: Current to voltage converting resistance calibration

Chapter 7.3.

Piezo actuators

We anticipated that the orientation of the mirror mount would not remain constant over time. In order to align the mirror in vacuum, the mirror mount is also accommodated by a pair of piezo transducers, in addition to the mechanical actuators. These are bound to the mirror mount using a 2-component epoxy H77 from *Epoxy technology*. A separate external piezo driver takes two control voltages between 0 and 10 V, in order to apply -10 to $+150$ V on individual piezos. As the accuracy and stability of the piezo elements affect the retro-reflection alignment, the following section describes the characterisation of its repeatability and thermal stability.

4.4 Parallelism alignment

The interferometer is sensitive to acceleration along the effective wave vector, $\vec{k}_{\text{eff}}$, which is defined by the Raman beams. For a small angle θ of misalignment between the Raman beams, $\vec{k}_1$ and $\vec{k}_2$, the interferometer phase can be written as

$$\Delta\Phi \approx k_{\text{eff}} \cos(\theta/2) a T^2. \quad (4.8)$$

This also affects the period of the interference fringes. To achieve a measurement accuracy of $\Delta a/a = 10^{-7}$, we need to reach sub-mrad alignment precision in parallelism, with $\theta \approx 0.9$ mrad.

In our setup, we exceed this requirement by retro-reflecting the Raman beams back into the collimator and by coupling the light back into the vacuum fibre. Figure 4.14 shows the setup for achieving retro-reflection. A 99-to-1 fibre splitter is used to guide the laser beam into the chamber. The 1%-port of the splitter is used to monitor the

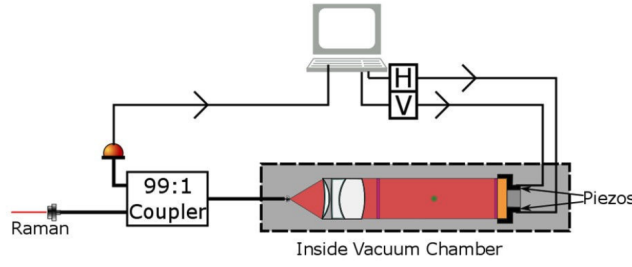


Figure 4.14: Schematics of the setup for the retro-reflection alignment: the 1% port of a 99-to-1 fibre splitter is used to monitor the power of the retro-reflected beam. The horizontal and vertical piezo actuators maximise the efficiency of the coupling.

retro-reflected light. Piezo transducers are controlled from a computer to maximise the optical power at the 1%-port. Once the back-coupling is successful, we can assure that the misalignment angle is reduced to $\leq 17 \mu\text{rad}$ ⁵. To maintain the alignment, the stability of the piezo transducers needs to be characterised.

Piezo characterisation

The stability of the piezo transducers is monitored by recording the position of a collimated laser beam, which is reflected from the piezo driven mirror mount. A well-calibrated screen is setup at 10.5 m away from the mirror. We have measured a maximal deflection angle of the light to be 1.2 mrad by applying a supply voltage between -10 and 150 V for each piezo transducer on the mirror. The output voltage of the piezo driver fluctuates by ± 0.1 V (peak-to-peak) over a measurement of 10 s, corresponding to $\pm 0.8 \mu\text{rad}$, which is negligible for us.

Figure 4.15 shows the behaviour of the piezo transducers as a function of temperature between 20° and 50°C . For each temperature, we measure the mean and the standard deviation of the deflection angle from 10 measurements. The angles of deflection are plotted against temperature in (a) for both piezo transducers in red and blue, respectively. Linear fits (solid lines) yield a temperature dependence of around $30 \mu\text{rad}/^\circ\text{C}$ for both elements. Because the laboratory temperature is stable to better than 1 degree, we expect the angle of deflection to stay well within the desired limit of 1 mrad. Figure 4.15 (b) shows the standard deviation of the deflection angle as a function of temperature. Both piezos exhibit a very satisfactory level of reproducibility. However, the angle is not stable enough to ensure retro-reflection back into the fibre. To compensate for the drift in the mirror angle, we apply a fast feedback algorithm, so-called the *Simplex algorithm*, which

⁵The fibre core has a diameter of approximately $5 \mu\text{m}$. As the ratio of the focal lengths of the aspherical lens pair is approximated to be 3, the focal point of the retro-reflected light after the triplet lenses need to be within $1.7 \mu\text{m}$ to the ideal focal point of the aspherical lens pair, in order to successfully couple the light back into the fibre. The mirror is positioned around 10 cm away from the collimator, which translates to the maximal misalignment angle of the light of $17 \mu\text{rad}$.

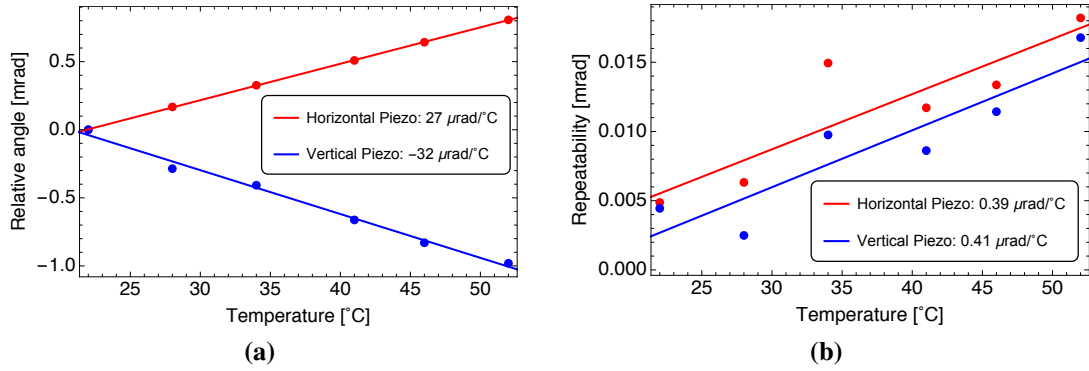


Figure 4.15: Temperature dependence of the piezo transducers. Relative change of the deflection angles of the laser beam is shown in (a), whereas the repeatability of each measurement is shown in (b) as a function of temperature.

automatically maximises the light coupled back into the vacuum fibre.

4.4.1 Automation alignment: Simplex algorithm

The simplex algorithm, proposed in Ref. [68], applies to multi-dimensional optimisation problems. Here, we wish to maximise the power reflected back in to the fibre as an objective function of two control parameters - the horizontal and vertical piezo control voltages. The basic idea is to start with any three points on the surface of the objective function, which form the initial geometric simplex. The algorithm iteratively performs reflection, contraction and expansion on the simplex according to the decision tree, illustrated in Fig. 4.16, until the power coupled into the fibre (or the coupling efficiency) converges to a maximum. The initial three points are first sorted in ascending order by the values of the coupling efficiency, and are denoted by L , M and H , as shown in (a). To replace the point L having the lowest coupling efficiency, the simplex is then reflected about the line between H and M . Following the decision tree in (b), a new measurement at the point R is compared with H , M and L , which then determines if the simplex needs to be expanded or contracted to a new point (E , C_e , C_i , C_{er} or C_{ir}). Once a new simplex is formed, the algorithm repeats.

The coloured contour plot in Fig. 4.17 (a) shows the measured coupling efficiency as a function of the horizontal and vertical piezo control voltages. The red dashed line indicates a real voltage changes during the optimisation with 27 iterations. I have also simulated 1000 such optimisations in order to assess how many steps are required to reach 50% of the optimum efficiency. The histogram in (b) summarises the results. Over 50% of the simulations have reached the goal within 15 iterations.

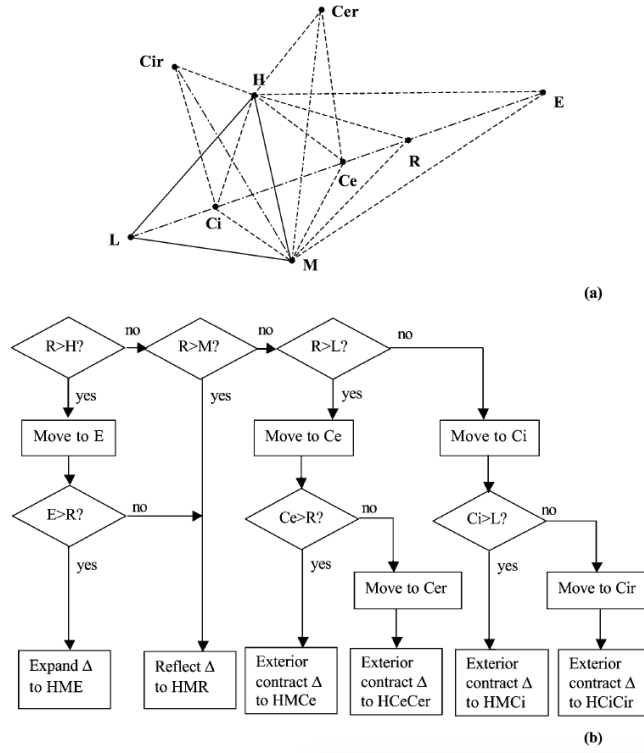


Figure 4.16: Simplex algorithm: (a) represents vectorially the parameters in a 2-dimensional optimisation problem. The decision tree for a single iteration is sketched in (b). Figures taken from Ref. [68].

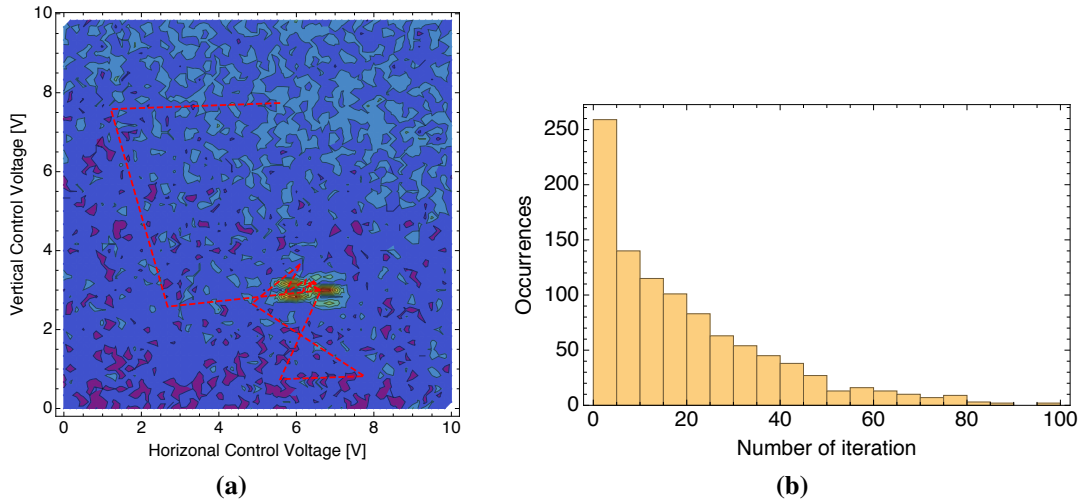


Figure 4.17: Experimental realisation of the simplex algorithm (a): The contour plot exhibits a map of measured coupling efficiency as a function of horizontal and vertical control voltages, where blue highlights a low coupling efficiency and red for a high efficiency. The red dashed line indicates 27 iterations of the algorithm before reaching 50% of the optimal retro-reflection. (b) summarises the results of 1000 simulated optimisations on the data map in (a). 50% of simulations occur to have reached the goal after 15 iterations.

Chapter 5

Raman Interaction

Our Raman optics setup, described in the last chapter, is designed to drive the velocity-sensitive σ^+/σ^+ or σ^-/σ^- clock transition $|F = 1, m_F = 0\rangle \leftrightarrow |F = 2, m_F = 0\rangle$. In this chapter, I describe and analyse the experimental observations after applying single and multiple double pulses.

5.1 Raman pulse

The transition probability of a ^{87}Rb atom after a two-photon stimulated Raman transition is described by Eq. 2.8. However, for the transition to be efficient, W in Eq. 2.9 needs to approach the effective Rabi frequency Ω_{eff} . This happens on resonance, when $\delta_{\text{AC}} = 0$ and $\delta = 0$. For vanishing AC stark shift, to resonantly address Raman transitions in moving atoms, the microwave beat frequency Δf_L between the Raman lasers needs to be on resonance with the clock transition at frequency, $f_0 = 6.834682 \text{ GHz}$, plus the Doppler shift $f_D \simeq 2v/\lambda$:

$$\Delta f_L \simeq f_0 + 2v/\lambda, \quad (5.1)$$

where v is the velocity of the atoms along the direction of $\vec{k}_{\text{eff}}$.

After the first three steps of the state preparation (see section 3.8), approximately 10^8 ^{87}Rb atoms are cooled down to few μK and prepared in the $|1, 0\rangle$ state. In an atom cloud at temperature T , the probability distribution for f_D is given by the Maxwell-Boltzmann distribution

$$P_{\text{MB}}(f_D) = \sqrt{\frac{m}{2\pi k_B T}} \exp\left(-\frac{m \lambda^2 f_D^2}{8 k_B T}\right). \quad (5.2)$$

The red dashed line in Fig. 5.1 (a) illustrates this distribution for a $8 \mu\text{K}$ cloud, which has a FWHM approximately 167 kHz . When a Raman pulse with a pulse duration of τ (fulfilling $\Omega_{\text{eff}}\tau = \pi$) is applied to an atom having no Doppler shift, the transition probability is given by Eq. 2.8, a sinc^2 function, whose width can be approximated by

$\Delta f_{\text{FWHM}} \approx 0.886/\tau$. The green and the black filled curves in Fig. 5.1 (a) show this probability as a function of the driving frequency for $\tau = 10 \mu\text{s}$ and $\tau = 100 \mu\text{s}$. They have FWHM of 88.6 kHz and 8.9 kHz, respectively. The overlap between the distribution of Doppler shifts and the line shape of Raman pulse gives an indication of which velocities are most effectively addressed by the Raman pulse. The long Raman pulse transfers only

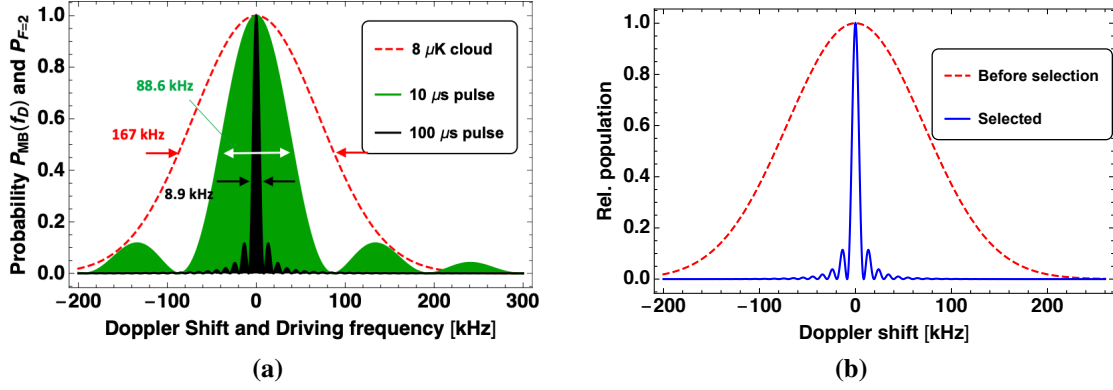


Figure 5.1: Probability distributions. Dashed red lines: distribution of Doppler shifts in an 8 μK cloud. (a) Shaded regions: transition probability versus driving frequency for an atom at rest, using rectangular Raman pulses of 10 μs (green) and 100 μs (black) duration. (b) Blue line: distribution of Doppler shifts for atoms driven to the state $|2, 0\rangle$ by a 100 μs Raman pulse tuned to f_0 .

a narrow group of velocities from one ground state into another one, and we are therefore able to use that as a velocity selection pulse. In Fig. 5.1 (b), the red dashed line (again) shows the distribution of Doppler shifts in the initial $|1, 0\rangle$ population. The blue line shows the probability of an atom being driven to the state $|2, 0\rangle$ by a 100 μs Raman pulse, as a function of its Doppler shift. This is calculated by weighting Eq. 2.8 (including the Doppler shift) with the distribution P_{MB} . Figure 5.2 (a) illustrates the fraction of atoms driven from the state $|1, 0\rangle$ to the state $|2, 0\rangle$ as a function of Raman pulse duration. The three curves are for clouds of temperature 8, 15 and 30 μK . This shows that a shorter

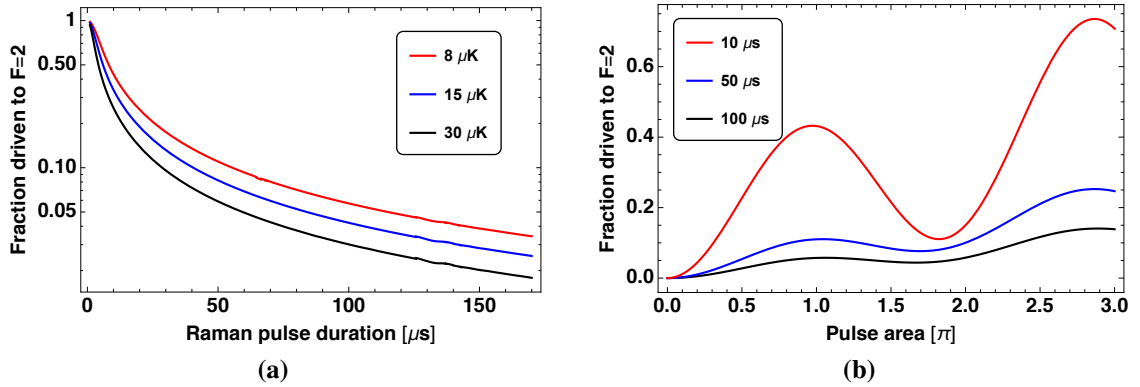


Figure 5.2: Fraction of atoms driven to $F = 2$: (a) as a function of Raman pulse duration for three different cloud temperatures. (b) as a function of the pulse area for three different pulse durations.

Raman pulse selects more atoms due to a broader frequency width of the selection pulse. At the same time, more atoms can be selected in a colder cloud. Figure. 5.2 (b) shows the fraction of atoms being driven to $F = 2$ as a function of the pulse area $\Omega_{\text{eff}} \tau$. The three curves are for pulses of duration for a $\tau = 10, 50$, and $100 \mu\text{s}$. All three curves exhibit Rabi oscillation, but this is most pronounced for the shortest pulse, because that addresses the widest group of velocities. For longer pulses, convolution with the distribution of Doppler shifts has the potential to spoil the visibility of the interferometer fringes. For this reason, we first use a long Raman pulse to pre-select a narrow group of velocities, for which the Doppler width is small in comparison with the width $1/\tau$ associated with the interferometer pulses¹.

In the remainder of this section, I describe how I setup the velocity selection pulse and characterised the resulting velocity distribution.

Single Raman pulse: velocity selection pulse

We start with all the atoms in the $|1, 0\rangle$ state. A single Raman pulse then transfers atoms into $F = 2$, and these are detected by resonant fluorescence (see section 3.8.1). Figure 5.3 (a) show the final $F = 2$ population of a $8 \mu\text{K}$ cloud as a function of the input Raman pulse power measured outside the vacuum chamber. The laser beat frequency is set at 6.8345 GHz to be resonant with the central Doppler-shifted frequency of the launched cloud (see section 5.2). The red and blue dots are the experimental data collected using Raman pulse durations of $\tau_{\text{vel}} = 40$ and $80 \mu\text{s}$, respectively.

To simulate the experiment, I average the $F = 2$ probability over the whole cloud, $P_{F=2}^{\text{vel}}$, as a function of Raman laser power p_L measured outside the chamber. This is described by

$$P_{F=2}^{\text{vel}}(p_L) = \int_{-\infty}^{\infty} P_{\text{MB}}(v, T) \cdot P_2(v, \eta_L p_L, \tau_{\text{vel}}) dv + \gamma_{\text{op}} \tau_{\text{vel}} \eta_L p_L. \quad (5.3)$$

The velocity distribution of the cloud at a temperature T is given by $P_{\text{MB}}(v, T)$. The transition probability $P_2(v, \eta_L p_L, \tau_{\text{vel}})$ is given by Eq. 2.8 after including the Doppler shift. η_L is the transmission of the Raman laser beam into the vacuum chamber. I assume that the intensity at the site of the atoms is that at the centre of a Gaussian beam with power $\eta_L p_L$ and a spot size of 38 mm . The second part of the Eq. 5.3 describes the contribution of off-resonant optical pumping into $F = 2$ by the Raman slave laser. This can drive $|F = 1\rangle \leftrightarrow |F' = 1\rangle$ and $|F = 1\rangle \leftrightarrow |F' = 2\rangle$, then the excited atoms can

¹A common way to suppress the side lobes of the velocity distribution selected by a top-hat pulse is to use Blackman-shaped velocity-selection pulses. More sophisticated pulse shaping techniques have been demonstrated in Ref. [69] with significant improvement of the interferometer contrast.

spontaneously decay back into $|F = 2\rangle$. In Chapter 6, we measure this optical pumping rate and find that $\gamma_{\text{op}} \approx 0.0011/(\mu\text{s W})$. The dashed lines in Fig. 5.3 (a) show the optical pumping contribution to $P_{F=2}$ as a function of input Raman laser power.

The solid lines in Fig. 5.3 (a) show the prediction of Eq. 5.3 for a $8\ \mu\text{K}$ cloud with Raman pulse duration of $\tau_{\text{vel}} = 40$ and $80\ \mu\text{s}$. I have treated the transmission η_L as a free parameter². To match the calculation with the data, η_L is set to be approximately 0.3. The calculation agrees with the data for low power. At higher power, only a very weak Rabi oscillation could be observed. In trying to understand the discrepancy between experiment and theory in Fig. 5.3, I eventually found (see Chapter 6) that the intensity distribution of the Raman beam had an unexpected lump in it, which introduced a significant variation of the Rabi frequency across the atom cloud. We think that is the reason for the discrepancy.

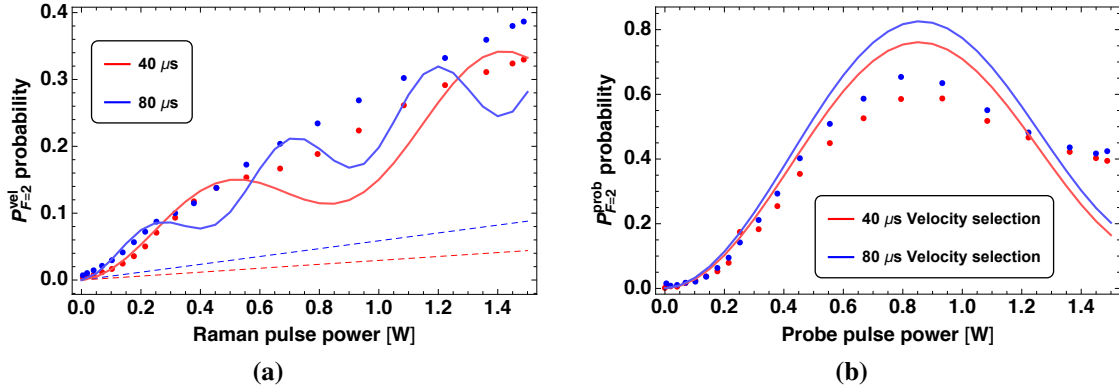


Figure 5.3: Rabi oscillations: (a) Raman pulse power scan for a single 40 and a $80\ \mu\text{s}$ pulse, compared with theoretical simulations (solid lines). (b) Power scan of a $20\ \mu\text{s}$ probe pulse after a 40 and a $80\ \mu\text{s}$ velocity-selective pulse.

Probing the velocity selection pulse

To characterise the velocities selected by the Raman pulse, we apply a second, short Raman pulse (referred to as a “probe” pulse) to the velocity-selected atoms. If the spread of Doppler shift of the velocity-selected atoms is much narrower than the spectral width of the probe pulse, the Rabi oscillation induced by the probe pulse should approach full visibility. To observe this, a narrower velocity class of atoms is selected from $F = 2$ into $F = 1$ using one long, velocity-selective Raman pulse. The selected atoms are then driven back into $F = 2$ by a $\tau_{\text{prob}} = 20\ \mu\text{s}$ probe pulse. Figure 5.3 (b) shows the probability of atoms on $F = 2$ after the probe pulse as a function of the input power of the probe pulse. Red and blue dots again present the measured data with velocity

²As the Raman optics are mounted inside vacuum, we could not measure the power transmission through the vacuum fibre feedthrough in-situ.

selection pulses optimised at $\tau_{\text{vel}} = 40$ and $80 \mu\text{s}$ respectively. The solid lines show the calculations of the $F = 2$ population after the probe pulse, $P_{F=2}^{\text{prob}}$, which are carried out using Eq. 5.4:

$$P_{F=2}^{\text{prob}}(p_L) = \int_{-\infty}^{\infty} \rho_{\text{vel}}(T, v, \eta_L p_L, \tau_{\text{vel}}) \cdot P_2(v, \eta_L p_L, \tau_{\text{prob}}) dv + \gamma_{\text{op}} \tau_{\text{prob}} \eta_L p_L. \quad (5.4)$$

Here, ρ_{vel} is the normalised distribution of Doppler shifts of the velocity-selected atoms and is described by

$$\rho_{\text{vel}}(T, v, \eta_L p_L, \tau_{\text{vel}}) = \frac{P_{\text{MB}}(v, T) \cdot P_2(v, \eta_L p_L, \tau_{\text{vel}})}{\int_{-\infty}^{\infty} P_{\text{MB}}(v', T) \cdot P_2(v', \eta_L p_L, \tau_{\text{vel}}) dv'}. \quad (5.5)$$

After comparing the data with the simulations, we conclude that a narrower velocity class significantly improves the Rabi oscillations, in contrast to Fig. 5.3 (a). However, the data shows a lower maximum transition probability and more dephasing than the simulations (solid lines), which I also attribute to the unexpected inhomogeneity of the intensity distribution in the Raman beams.

5.2 Raman spectrum

The narrow line width of the velocity selection pulse also enables us to study the velocity distribution of the entire atom cloud. To perform a wide scan of the distribution of Doppler shifts, we prepare atoms initially in $|1, 0\rangle$, and a Raman probe pulse transfers them into $F = 2$, which is measured by $P_{F=2}$ probability. Figure 5.4 shows $P_{F=2}$ as a function of Raman beat note frequency, Δf_L , optimised for maximum peak transition probability at four different pulse durations τ . The resolution of the distribution is determined by the frequency width of the Raman probe pulse (i.e. $\approx 0.886/\tau$). Three resonances can be seen in Fig. 5.4. As expected, Doppler-broadened lines appear symmetrically around the frequency of the atomic hyperfine transition $|1, 0\rangle \leftrightarrow |2, 0\rangle$ at $f_0 = 6834.682 \text{ MHz}$ [57]. This is due to the presence of two separate velocity-sensitive Raman transitions with opposite momentum transfer ($\pm \hbar k_{\text{eff}}$) (see section 2.1.3). For atoms moving at a velocity of approximately 7 cm/s along the Raman beam axis, the Doppler shift is upward for one pair of counter-propagating laser beams and downward for the other, which cause these Doppler-shifted velocity-sensitive peaks. In the centre of the spectrum, there is a narrow, velocity-insensitive peak, which is due to excitation by co-propagating beams and indicates a misalignment of the QWP. The width of this line can be used to check the duration of the probe pulse, while its centre frequency gives a Doppler-free measure of the clock transition frequency.

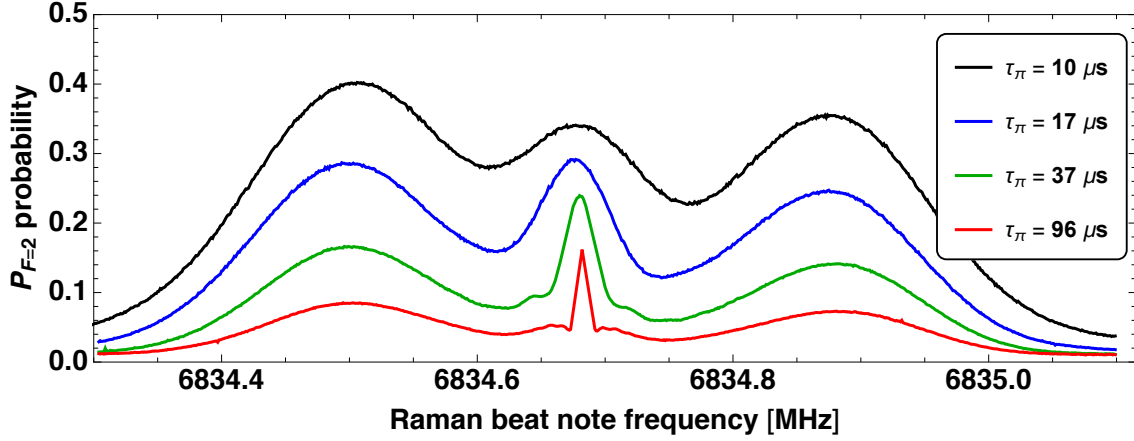


Figure 5.4: Raman spectra probed by a $\tau_\pi = 10, 17, 37$ and $96 \mu\text{s}$ Raman pulse: Atoms are prepared initially in $|1, 0\rangle$ and detected with $F = 2$ fluorescence detection. The central peaks indicate polarisation misalignment in the Raman beams.

Temperature estimate

The Doppler-broadened lines in Fig. 5.4 can be modelled by convolving the Doppler-shifted version of Eq. 2.8 with the underlying velocity distribution of the atom cloud. If the spectral width of the Raman probe pulse is negligible with respect to the range of Doppler shifts, we can interpret the measured line shape to be the underlying velocity distribution. Now, the data in Fig. 5.4 are partitioned into datasets that only include velocity-sensitive lines, and I assume that the velocity distribution of the cloud is a Maxwell-Boltzmann. Gaussian fits ($\propto \exp(-\frac{f^2}{2(\Delta f_{\text{con}})^2})$) determine the standard deviation Δf_{con} and hence define a temperature T given by

$$T = \frac{m_{\text{Rb}} \lambda^2 \Delta f_{\text{con}}^2}{4 k_B}. \quad (5.6)$$

The data points in Fig. 5.5 (a) show the temperature determined in this way for a variety of probe pulse durations τ . Because of the convolution, we see that the temperature of the atoms is inflated by this method when τ is short. In order to make a simple approximate correction for this effect, I take the frequency spectrum of the Raman probe pulse to be a Gaussian with a standard deviation of approximately $0.38/\tau$, which approximates the line width of the rectangular Raman probe pulse with a pulse duration of τ . The measured width Δf_{con} can then be expressed by

$$\Delta f_{\text{con}} = \sqrt{\Delta f_{\text{atom}}^2 + \left(\frac{0.38}{\tau}\right)^2}, \quad (5.7)$$

where Δf_{atom} gives the actual velocity spread of the atom cloud. The red dashed line in Fig. 5.5 (a) illustrates the fit of the temperature using Eq. 5.7, which gives a cloud tem-

perature of approximately 7 μK . This results is in good agreement with the temperature obtained from the time-of-flight method discussed in section 3.7.2.

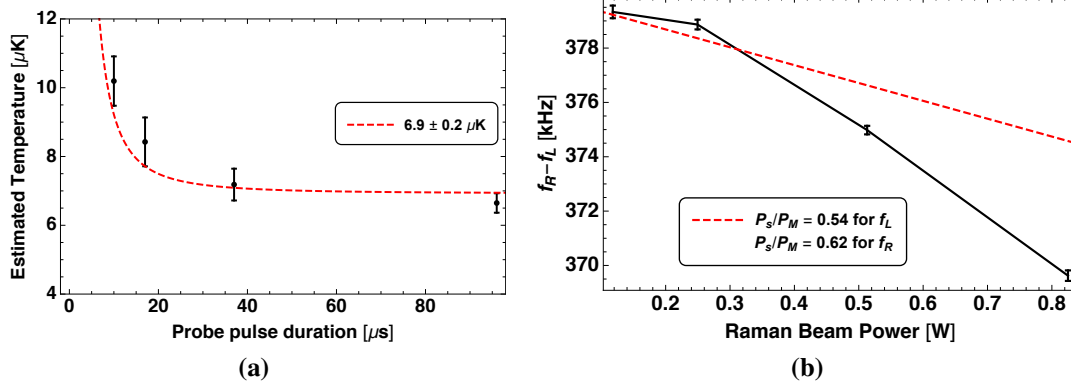


Figure 5.5: (a) Molasses "temperature" as given by the squared width of the velocity-sensitive Raman spectrum. On correcting for the effect of probe pulse duration, the fit gives a temperature for the atoms of 6.9 μK , which agrees with the time-of-flight measurement. (b) the differential light shift as a function of the relative Raman laser power. The red dashed line illustrates a simulation of the light shift with imperfect mirror reflectivity, where the input intensity ratio is set to 0.58 and the horizontal launch velocity of 190 kHz.

Remarks on the light shift

The light shift can be measured from the shift of the central frequency of each envelope in Fig. 5.4. Let us write the central frequencies of the two Doppler-broadened lines as f_L and f_R . These can be described in terms of absolute frequencies by

$$\begin{aligned} f_L &= f_0 - \Delta f_H + \Delta f_L^{AC} + \Delta f_{\text{Zeeman}}, \\ f_R &= f_0 + \Delta f_H + \Delta f_R^{AC} + \Delta f_{\text{Zeeman}}, \end{aligned} \quad (5.8)$$

where f_0 is the unperturbed hyperfine transition frequency, Δf_H is the Doppler shift due to the horizontal launch velocity of the atom cloud, Δf^{AC} specifies the light shift caused by the Raman beams, and Δf_{Zeeman} is the second order Zeeman shift of the clock transition (575 Hz/G²). In the applied field of 1.4 G, this Zeeman shift amounts to approximately 1 kHz. The splitting between the two lines is sensitive only to the Doppler shift produced by launching and the difference of the light shifts coming from the two pairs of counter-propagating Raman beams, associated with $\pm \hbar k_{\text{eff}}$ momentum transfer:

$$f_R - f_L = 2\Delta f_H + (\Delta f_R^{AC} - \Delta f_L^{AC}). \quad (5.9)$$

If the laser beam pairs have the same intensities, Eq. 5.9 becomes independent of the Raman laser power, because $\Delta f_R^{AC} = \Delta f_L^{AC}$. However, the retro-reflecting mirror has

only 93% reflectivity, which changes the laser intensity ratio seen by the atoms to 0.54 and 0.62 separately for the Raman beam pairs with $\pm \hbar k_{\text{eff}}$ momentum transfer. Hence, we expect a linear variation of the differential central frequency, $f_R - f_L$ in Eq. 5.9, between the Doppler-shifted distributions due to the light shifts, as a function of the probe pulse power. The black connected points in Fig. 5.5 (b) shows Eq. 5.9, as determined in this way from the Raman spectrum in Fig. 5.4. The red dashed line shows the calculation of $f_R - f_L$ ³. By assuming pure σ^+/σ^+ transitions, the light shifts, $\Delta f_{L/R}^{AC}$, are estimated by summing up all contributions to the clock states due to the coupling of the Raman beams with all accessible states on the D1 and D2 lines (including the expected intensity ratios of 0.54 and 0.62). The nonlinearity evident in Fig. 5.5 (b) is probably because the fibre-to-fibre connector going through the vacuum wall affects the polarisation and the power transmitted in a given polarisation, in a way that depends on the power.

It is worth mentioning that the offset shift ($2\Delta f_H \approx 380$ kHz) in Fig. 5.5 (b) corresponds to an initial launch velocity of the cloud being 7.4 cm/s, which differs by 0.6 cm/s from the velocity estimated using camera as described in section 3.7.1. This could be explained by a tilting angle of 2° of the optical table away from the horizontal. In that case, the component of gravity along the Raman axis would deaccelerate the atoms, causing them to lose 0.62 cm/s of velocity at the time of the measurement (15.5 ms after the launch).

A more complex behaviour is observed with frequency shift of the velocity-insensitive transition peak, which is not presented in the thesis and requires further investigation. To minimise the polarisation instability of the Raman beams in the next version of the setup, a single-mode polarisation-maintaining fibre will be used to guide the laser light directly through the vacuum wall without any connector.

5.2.1 Velocity selection pulse

As seen in the section 5.1, atoms with a narrow velocity distribution exhibit a significantly better Rabi oscillation. It is interesting to analyse the Raman spectrum after the velocity selection pulse.

The black dots in Fig. 5.6 shows the upper Doppler-shifted distribution of the atoms around the Raman beat note frequency of 6834.89 MHz, which is driven by a 80 μs Raman pulse. This shows the Raman spectrum before the velocity selection pulse and is the same measurement as described in Fig. 5.4. The FWHM of the distribution is estimated to be approximately 188 kHz. To obtain the Raman spectrum of the velocity-selected atoms, we prepare atoms initially in the state $|2, 0\rangle$. A velocity selection pulse with $\tau_{\text{vel}} = 40 \mu\text{s}$

³The credit for the calculations of the light shifts goes to Prof. Hinds.

is applied at a Raman beat note frequency of 6834.89 MHz. This is resonant with the central frequency of the upper Doppler-shifted velocity-sensitive line driven by the Raman beam pairs with $+\hbar k_{\text{eff}}$ momentum transfer. After some atoms are selected into $|1, 0\rangle$, the atoms remaining in $F = 2$ are cleared out using a resonant blow-away beam. Next, a $80\text{ }\mu\text{s}$ Raman probe pulse drives the selected atoms back to $F = 2$, where they are detected by resonance fluorescence. This is shown in $P_{F=2}$ probability as a function of the Raman beat note frequency of the probe pulse in blue line in Fig. 5.6 (a).

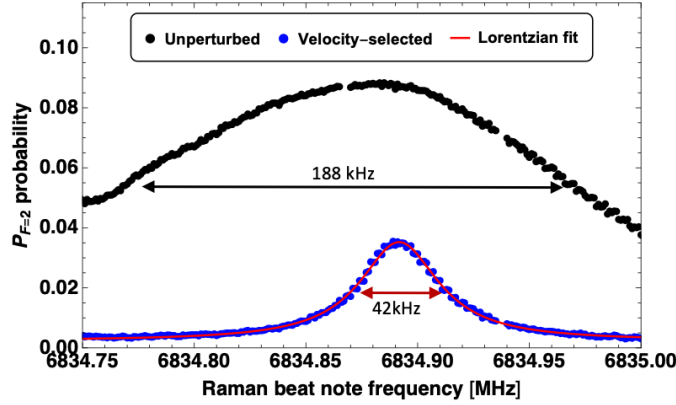


Figure 5.6: Raman spectrum: Raman spectrum of unperturbed atoms before the velocity selection pulse (black dots), which is driven by a $80\text{ }\mu\text{s}$ Raman pulse. Atoms selected by a $40\text{ }\mu\text{s}$ velocity selection pulse are probed by a $80\text{ }\mu\text{s}$ Raman pulse, shown in blue dots. The solid red line exhibits the Lorentzian fit of the data in blue. Here, the velocity selection pulse manages to reduce the line width of velocity distribution from 188 kHz to 42 kHz at FWHM.

It is worth noting that the expected line width of the peak at 6834.89 MHz in blue in Fig. 5.6 (a) is around 25 kHz by convolving two sinc^2 functions, where one of them describes the frequency distribution of the $40\text{ }\mu\text{s}$ velocity selection pulse and the other one is the distribution of the $80\text{ }\mu\text{s}$ Raman probe pulse⁴. However, a Lorentzian fit (red line) in Fig. 5.6 (b) exhibits a width of 42 kHz. This discrepancy could be explained by imperfect Raman pulses. The frequency distribution of an imperfect Raman pulse can no longer be described by a perfect sinc^2 function, as the side lobes in the sinc^2 function gradually gain their significance, which result in the line broadening of the Raman pulse. I have simulated this broadening effect and could reproduce the observed width of 42 kHz after convolving two imperfect Raman pulses. Imperfect Raman pulses could be caused by the intensity inhomogeneity of the Raman beam cross the atom cloud (see Chapter 6). Despite of this, the velocity distribution of atoms is reduced by at least a factor of 4 after the velocity selection pulse, which is sufficient for observing interference fringes.

⁴This can also be very well approximated by convolution between two Gaussians

Chapter 6

Fringe optimisation

Reference [62] focused on making first observations of the interference fringes. The focus of my research was to improve the fringe contrast, which, in turn, decreased the acceleration uncertainty σ_a of the interferometer. In the last chapter, I described the use of a long Raman pulse to select atoms in a narrow range of velocities along $\vec{k}_{\text{eff}}$. The purpose of this was to ensure that the pulse areas of the interferometer pulses are not compromised by the velocity spread. In this chapter, I describe my progress on the interferometer itself.

6.1 The first interference fringes

We start by setting the pulse areas of the interferometer pulses. With the Raman laser at maximum output power we vary the pulse duration¹, and we take the π -pulse time, τ_π , to be when the maximum in the $F = 2$ population is reached the first time. Similarly, we take the $\pi/2$ -pulse time, $\tau_{\pi/2}$, to be when the $F = 2$ population has reached half of the population of a π -pulse. Next, we scan the beat-note phase of the π -pulse, ϕ_2 , to observe interference fringes

Figure 6.1 presents the first interference fringe observed using our 3-pulse interferometer with a $2T$ -time of 50 ms. The $F = 2$ probability of atoms after the interferometer is plotted as a function of ϕ_2 . We use Eq. 6.1,

$$P_{F=2} = P_0 + \frac{C}{2} \cos(2\phi_2 + \phi_0), \quad (6.1)$$

to fit the fringe data, as shown in red dashed line. P_0 is the offset signal and ϕ_0 is the offset interferometer phase. The fringe contrast C is only 0.046 ± 0.003 .

¹It is perhaps more natural to vary the pulse power while keeping the pulse duration constant. This avoids changes in the frequency width of the scanning pulse. As the Doppler width of the velocity-selected atoms is much narrower than the spectral width of the scanned pulse, we didn't observe any noticeable difference between the two methods.

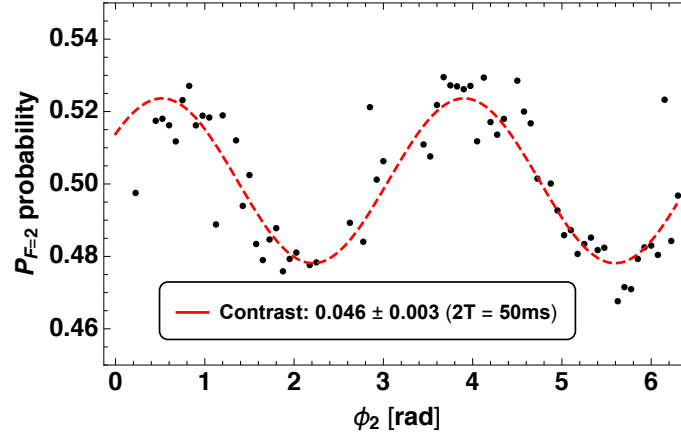


Figure 6.1: Interference fringe scan for an interferometer with $2T = 50$ ms. A velocity-selection pulse of $80 \mu\text{s}$ is applied, in combination with two $13 \mu\text{s}$ $\pi/2$ -pulses and a $30 \mu\text{s}$ π -pulse. Data has been corrected by removing the contribution of $|1, \pm 1\rangle$ population.

There are several factors, which can cause such low fringe contrast. A reduction of the contrast can indicate decoherence in the atomic system via any external excitations or collisions with background atoms [70]. Such interactions destroy the coherent superposition in an atom, when the atom reveals its which-way-information during the interferometry. As external laser and microwave sources are switched off between the interferometer pulses, it is unlikely for any excitation to take place. To suppress collisional decoherence, the number of collisions between a cold Rb atom and the background residual gases during the interferometry needs to be $\ll 1$. For simplicity, I assume that the background residual gas is mainly hydrogen. I approximate the number of collisions by $\sigma_{\text{Rb-H}_2} n \bar{v}_{\text{H}_2} 2T$, where the total collisional cross-section between Rb and H_2 is $\sigma_{\text{Rb-H}_2} \approx 3 \times 10^{-14} \text{ cm}^2$ [71], n is the mean number density of H_2 at room temperature, and $\bar{v}_{\text{H}_2} \approx 1700 \text{ m/s}$ is the mean velocity of H_2 molecules. Figure 6.2 shows the estimated number of collisions during a $2T = 50$ ms interferometer as a function of the background hydrogen pressure. In UHV with $< 5 \cdot 10^{-10} \text{ mbar}$, there is a chance of $< 1\%$ for a Rb atom to collide with a background hydrogen molecule. Hence, our vacuum pressure should not affect the fringe contrast. This leads me to investigate the reduction of contrast due to averaging over single-atom fringes across inhomogeneous intensity profile of the Raman beam.

6.1.1 Contrast characterisation

A variation of the Raman pulse areas corresponds to a variation of the Rabi frequencies across the atom cloud, which arises both from the Doppler width of the velocity-selected atoms and from the intensity variation of the Raman beam across the atom cloud. To analyse these contributions, we look at the interference fringe contrast as a function of

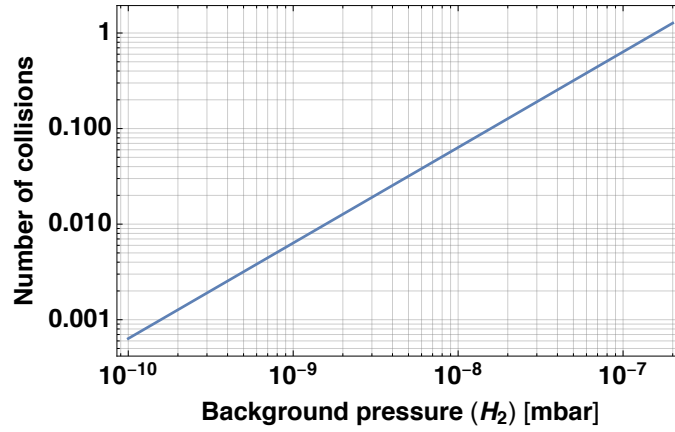


Figure 6.2: Estimated number of collisions between a ^{87}Rb atom and hydrogen molecules in the background gas within an interferometer with $2T = 50$ ms.

the interferometer time $2T$ for two different pulse durations, τ_{vel} , of the velocity selection pulse, as shown in Fig. 6.3. The connected blue and red data points show the fringe

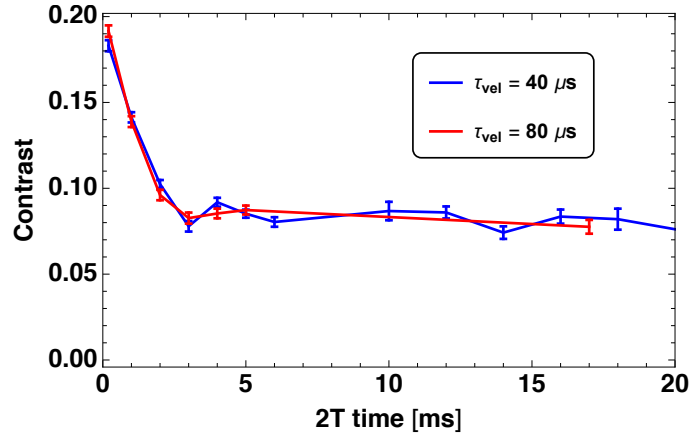


Figure 6.3: Interference fringe contrast plotted as a function of interferometer time $2T$ with $14\ \mu\text{s}$ of the first and third interferometer pulse and $28\ \mu\text{s}$ of the second interferometer pulse. A 40 and a $80\ \mu\text{s}$ velocity-selection pulse are applied at a detuning of -230 kHz from the atomic hyperfine transition (blue and red lines). In both cases, the fringe contrast reaches a constant value of 0.1 after 3 ms of interferometer time.

contrasts with a $\tau_{\text{vel}} = 40\ \mu\text{s}$ and a $\tau_{\text{vel}} = 80\ \mu\text{s}$ velocity-selection pulse, detuned at -230 kHz from the atomic hyperfine transition. The atom cloud was launched vertically with a velocity of $v_V = 25.7$ cm/s and horizontally, along the Raman beam axis with a velocity of approximately $v_H = 10$ cm/s. Both fringe contrasts reach a constant value of 0.1 for interferometer times between $2T = 3$ ms and 20 ms. The perfect agreement between the two sets of data indicates that the observed loss of contrast is not governed by the velocity spread of the velocity-selected atoms. This finding leads us to investigate the intensity profile of the Raman beams.

6.2 Trajectory analysis - beam profiling

In order to explore the variation of the Raman beam intensity, we launch the atom cloud on two different trajectories, illustrated in Fig. 6.4, using the moving molasses techniques, described in section 3.7.1. In one case, we launch the atoms horizontally along the Raman beam with a velocity of 7.5 cm/s, denoted by H . In the other case, we add a vertical component to the launch with a velocity of 25.7 cm/s against the gravity, labeled by HV . The Rabi oscillations of atoms moving along these two trajectories allow us to study the

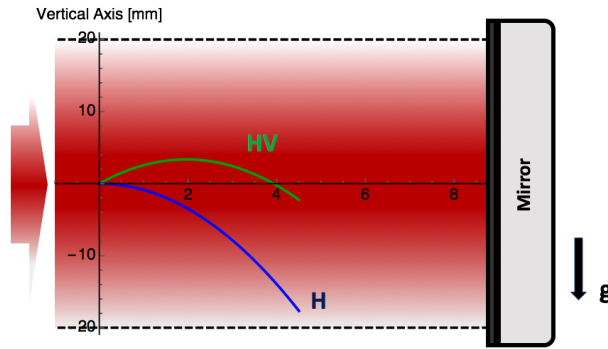


Figure 6.4: The horizontal and vertical launch trajectories is illustrated in green (HV), whereas the horizontal only launch trajectory is shown in blue (H).

distribution of intensity within the Raman beams. The effective Rabi frequency for the Raman transition (Eq. 2.10) is proportional to the product of the electric fields in the two counter-propagating Raman beams. For Raman transition via multiple intermediate levels, the effective Rabi frequency Ω_{eff} for 2-photon Raman transition is given by

$$\Omega_{\text{eff}} = \sum_j \frac{\Omega_{1,j}(E_1) \Omega_{2,j}^*(E_2)}{2 \Delta} \propto \sqrt{I_1 I_2} \propto I_0, \quad (6.2)$$

with $\Omega_{i,j}(E_i)$ being the Rabi frequency between $F = i$ ground state and the intermediate states j driven by the electric field E_i with the intensity I_i for $i = 1, 2$. Since these beams have a fixed ratio, η , between the master and the slave Raman laser, the intensity in each Raman beam can be expressed by $I_1 = I_0/(1 + \eta)$ and $I_2 = \eta I_1$ with I_0 being the total laser intensity. Therefore, I_0 can be deduced from Rabi oscillation. In order to observe the oscillation, we record $P_{F=2}$ as a function of pulse duration for a fixed laser power. To minimise the effect of the changing bandwidth of the probing pulse, a much narrower velocity class of atoms is initially selected by a 44 μs or 80 μs pulse. The probing pulse is set to maximum power, which provides us with a π -pulse time of around 12 μs , whose spectral width is significantly larger than the Doppler width of the velocity-selected atoms.

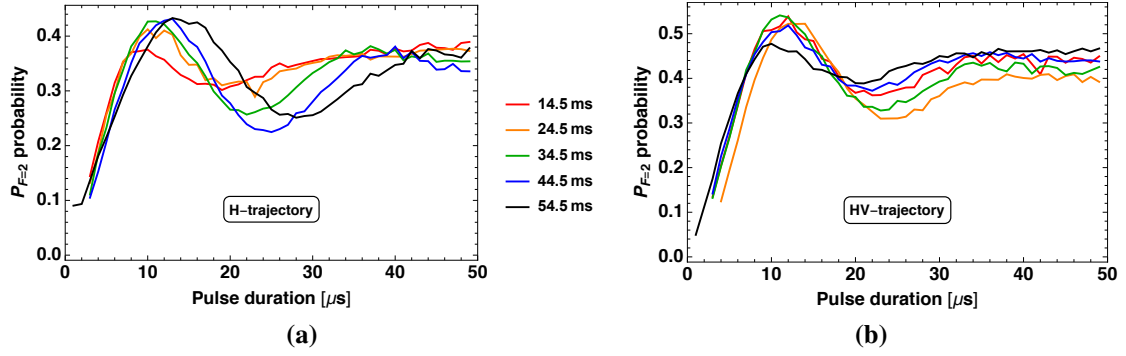


Figure 6.5: Rabi oscillations are performed on H -launched (a) and HV -launched (b) clouds for various flight times: Atoms are prepared in $|1, 0\rangle$ by a $44\ \mu\text{s}$ velocity-selection pulse. The duration of a probe pulse is scanned at the maximal laser power at $-180\ \text{kHz}$ red-detuned from the atomic hyperfine transition frequency.

Figure 6.5 shows the $F = 2$ probability as a function of the duration of the Raman probe pulse, which is applied after a $44\ \mu\text{s}$ velocity-selection pulse on atoms prepared in $|1, 0\rangle$. The coloured lines show the Rabi oscillations, which are collected at different flight times after the molasses being launched into H - (a) and HV -trajectories (b). Two features can be observed. As shown in both trajectories, the durations of π -pulse vary with the flight time of the molasses, which indicate a noticeable change of beam intensity in the direction of the launch trajectories. Secondly, the amplitudes of $P_{F=2}$ around the first maximum (π -pulse duration) and the first minimum (2π -pulse duration) both exhibit strong flight time-dependence, which, in turns, show position-dependent dephasing. In the following sections, I will use the variations of π - and $\pi/2$ -pulse durations and the corresponding $P_{F=2}$ probabilities to infer the intensity inhomogeneity of the Raman beam, and furthermore analyse the loss of fringe contrast. Firstly, let us look at the shift of the π -pulse durations.

6.2.1 π -pulse time measurement

To estimate the π -pulse duration, τ_π , where $P_{F=2}$ probability reaches the first maximum, we fit a parabola to the data points around the maximal probability in Fig. 6.5 (a) and (b). Using the initial launch velocities, we convert the flight time into the vertical position of the atoms, denoted as “height”. Figure 6.6 displays the measured π -pulse times as a function of the “height” during the H - and HV -trajectories (in blue and red respectively). Using $\tau_\pi = \pi/\Omega_{\text{eff}}$ and Eq. 6.2, the π -pulse time, τ_π , is proportional to I_0 ,

$$\frac{1}{\tau_\pi} \propto I_0. \quad (6.3)$$

Figure 6.7 exhibits the expected intensity profile for a Raman beam with a spot size

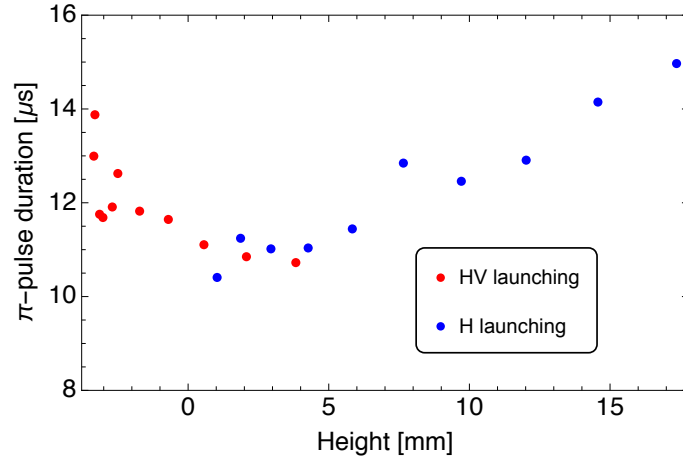


Figure 6.6: π -pulse times are estimated from data in Fig. 6.5 and plotted as a function of height, where the molasses is formed at 0 mm in height. Increasing height indicates the direction of the gravity.

of 37 mm in the dashed green line for comparison. The connected red dots, on the other hand, show the relative beam intensity as a function of the height, which is obtained from the data in Fig. 6.6 by using Eq. 6.3. We see that the inferred intensity varies less

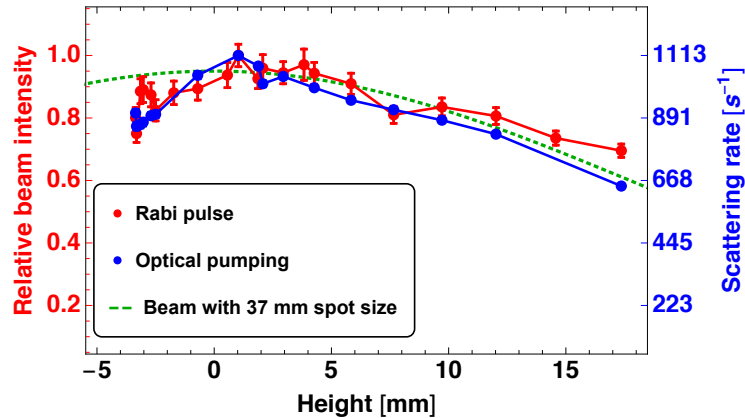


Figure 6.7: Intensity profiling: the relative beam intensity is measured by converting the π -pulse time into laser intensity, which are shown in red dots as a function of the height. The blue data points show the scattering rates estimated from optical pumping, which are measured with the master Raman laser solely. Both methods suspect the appearance of unexpected intensity lumps between ± 5 mm.

towards the right hand side of the graph and the variation seems to increase towards 0 mm. However, between 0 mm and -3 mm, there is a strong unexpected dip. The small error bars indicate that the intensity dip is perhaps a real physical feature in the beam, which requires further analysis. In the next section, I use optical pumping to confirm this finding.

6.2.2 Optical pumping

An alternative method of mapping the beam intensity profile is to use off-resonant, optical pumping. This is achieved by initially preparing all atoms in $F = 2$. When the master Raman laser solely is turned on with a frequency detuning of $\Delta = -1.13$ GHz from $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transition, the atoms are off-resonantly excited to the $5^2P_{3/2}$ states. Since the detuning is much larger than the line width of the transition Γ , the total scattering rate for the ground state ($F = 2$) can be approximated in the first order perturbation theory for laser field with the polarisation q by a sum over all possible hyperfine transitions, given by

$$R_{\text{sc}}^{\text{tot}, q} \approx \Gamma \sum_{F'=1,2,3} \frac{|\langle F = 2 | d_q E_q | F' \rangle|^2}{\hbar^2 \Delta^2}. \quad (6.4)$$

For 1 W of master laser power with around 50% coupling efficiency into the chamber, which is then expanded into a Gaussian with a spot size of 37 mm, we expect a total scattering rate of around 6400/s from $|2, 0\rangle$ using Eq. 6.4. To estimate the probability of an atom in $|2, 0\rangle$ being optically pumped to $F = 1$, we evaluate the branching ratio of the transition $|2, 0\rangle \rightarrow |F' = 2\rangle \rightarrow |F = 1\rangle$ being $\frac{1}{4} \times \frac{1}{2}$ and $\frac{1}{20} \times \frac{1}{2}$ for $|2, 0\rangle \rightarrow |F' = 1\rangle \rightarrow |F = 1\rangle$. After summing up the branching ratios, we expect a total optical pumping rate from $|F = 2, m_F = 0\rangle \rightarrow |F = 1\rangle$ of 960/s. As the scattering rate is proportional to the local beam intensity, the local optical pumping rate allows us to infer the intensity variation of the beam. In the experiment, we prepare atoms initially in $|2, 0\rangle$. Figure 6.8 shows the loss of these atoms in $P_{F=2}$ as we increase the pulse duration of the master laser at the maximum power of 1 W. Four different colours denote different flight times, which correspond to different heights. The solid lines in Fig. 6.8 show the linear fits, and

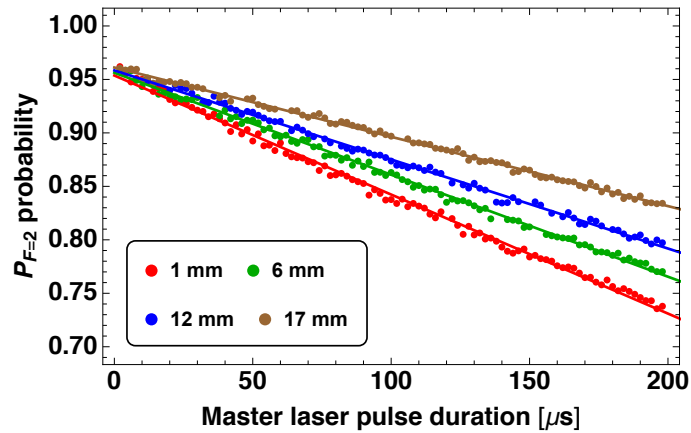


Figure 6.8: Optical pumping: Atoms are initially prepared in $|2, 0\rangle$. Using only the master laser, the loss rate of an atom being in $F = 2$ is measured as a function of the laser pulse duration. Different colours indicate different heights of the atom cloud in the Raman beam.

we determine the loss rate at shorter pulse duration from the slopes of these lines. The blue data points in Fig. 6.7 summarise the measured scattering rates as a function of the height. We see that optical pumping shows a similar intensity profile to the one obtained from π -time measurement, but with less noise. Again there is an unexpected intensity dip between -3 mm and 0 mm. Despite the dip, the intensity only varies by 10% to 20% around the centre of the beam, which is not enough to explain the low fringe contrast in Fig. 6.1 and the low visibility of the Rabi oscillations in Fig. 6.5. We wondered whether the intensity variation was actually greater, but at a higher spatial frequency that was unresolved across the 4 mm of the atom cloud. In the next section, I use the amplitudes of $P_{F=2}$ at π - and 2π -pulse times to measure the dephasing in the Rabi oscillations in order to test if there are indeed strong intensity variations of high spatial frequency.

6.2.3 Dephasing of the Rabi oscillation

The loss of visibility of the Rabi oscillation shown in Fig. 6.5 indicates dephasing of the Rabi oscillation. This is to be expected if there is indeed a significant intensity gradient across the size of the atom cloud. To interpret the data in Fig. 6.5, I assume that the beam intensity varies across the cloud linearly according to $I(x) = I_0 + \Delta I(x) = I_0 + \frac{dI}{dx} \cdot x$, where I_0 is the mean intensity seen by the atoms and x is the position within the cloud. This intensity gradient is incorporated into Eq. 5.4 with $P_{F=2}^{\text{prob}}(\tau_{\text{prob}}, I(x))$, and we obtain the overall $F = 2$ probability by averaging over the atom cloud using:

$$P_{F=2}^{\text{deph}}(\tau_{\text{prob}}) = \int_{-\infty}^{\infty} n_{\text{atom}}(x) P_{F=2}^{\text{prob}}(\tau_{\text{prob}}, I(x)) dx, \quad (6.5)$$

where $n_{\text{atom}}(x)$ is a Gaussian atom number density distribution. The black dots in Fig. 6.9 show the Rabi oscillation as a function of the probe pulse duration for atoms which are velocity-selected by a $80 \mu\text{s}$ velocity selection pulse. The measurement is performed at a height of 1.2 mm, which is around the intensity lumps, where the beam intensity varies by 10-20% over the size of the atom cloud according to Fig. 6.7. The red and blue dashed lines exhibit Eq. 6.5 for $\Delta I/I_0 = 0.1$ and 0.2 , respectively for a 10% and 20% intensity variation. The calculated Rabi oscillations for both cases are far more visible than those in the data. To resemble the data, a large intensity variation of 90% or $\Delta I/I_0 = 0.9$ is required, as shown by the dashed green line. Hence, we surmise that there are strong intensity variations that are unresolved in Fig. 6.7 but cause the damping of the Rabi oscillations seen in Fig. 6.5.

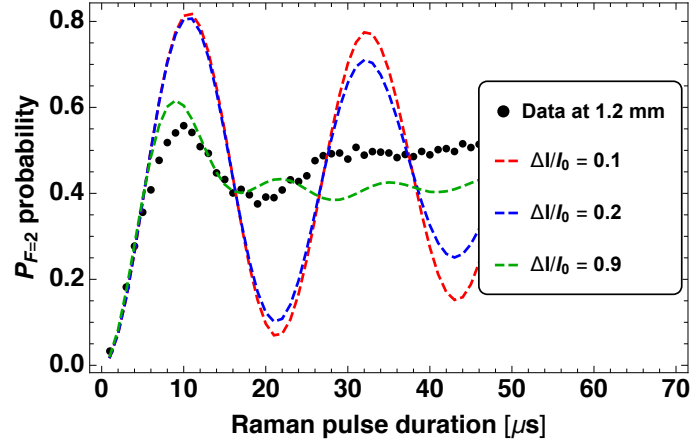


Figure 6.9: Rabi oscillations of atoms velocity-selected by a $80 \mu\text{s}$ Raman pulse as a function of the probe pulse duration. Data (black dots) are collected at a height of 1.2 mm within the intensity lump. The dashed lines show calculations of Rabi oscillations using Eq. 6.5 by assuming a linear intensity gradient of $\Delta I/I_0 = 0.1, 0.2$, and 0.9 .

To quantify this further, I define a Rabi flopping visibility

$$\eta_{\text{Rabi}} = \frac{P_{\text{max}} - P_{\text{min}}}{P_{\text{max}} + P_{\text{min}}}, \quad (6.6)$$

where P_{max} and P_{min} denote the maximum and minimum values of $P_{F=2}$, which correspond to the $F = 2$ probabilities at π - and 2π -pulse durations, respectively. A large

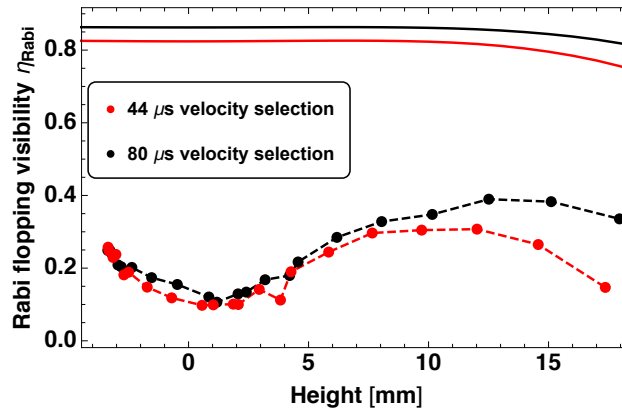


Figure 6.10: Rabi flopping visibility for a $8 \mu\text{K}$ atom cloud velocity-selected by a 44 (red) and $80 \mu\text{s}$ (black) Raman pulse as a function of height. A decrease is observed between $\pm 5 \text{ mm}$, which is congruent with the expected position of the intensity lump. The solid lines show the calculated Rabi flopping visibility for an ideal Gaussian Raman beam using Eq. 6.5.

intensity gradient decreases η_{Rabi} . The solid red and black lines in Fig. 6.10 show visibilities calculated using Eq. 6.5 for our ideal Gaussian intensity profile with a spot size of 37 mm and for our Gaussian atom cloud with 2 mm standard deviation. The two lines are for velocity selection by a 44 and a $80 \mu\text{s}$ Raman pulse. Both have high visibility, with the broader velocity class (red solid line) showing a slightly lower visibility because

it expands further towards the edge of the Raman beam where the gradient is stronger. The connected points in Fig. 6.10 show the measured visibilities as a function of height, derived from the individual Rabi oscillations in Fig. 6.5. Both sets of data exhibit low visibility, with a clear dip around 0 mm. At larger height, these data sets drift apart from each other as expected, but more than they do in the solid curves. More investigation is required to understand, why the velocity widths of the atom cloud contribute to the dephasing more significantly at the edge of the beam than predicted in the simulation.

This result supports our suspicion of strong intensity variations at high spatial frequency around the centre of the beam. It suggests that better interference fringe contrast could be achieved between 5 and 15 mm. To demonstrate this, I use a short 3-pulse interferometer to probe the beam profile in the next section.

6.2.4 Interferometric probing

To measure the fringe contrast at different heights, we use an interferometer with a $2T = 0.2$ ms. A short interferometer ensures that the atoms don't fall far during the interferometry, and hence it allows the intensity probing to be well localised. Further, the large fringe spacing² makes the interferometer insensitive to vibrations which would otherwise compromise the fringe contrast.

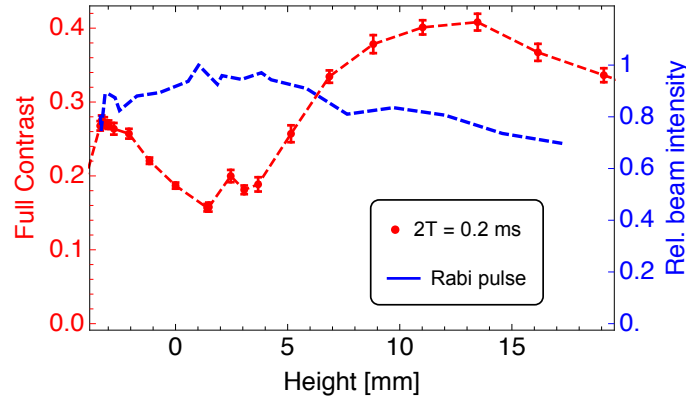


Figure 6.11: Comparison between the fringe contrasts (red) of a $2T = 0.2$ ms interferometer and the beam intensity measured by Rabi oscillations (blue).

Atoms are launched along the H and HV trajectories described previously. At each height, we vary the beat note phase of the second interferometer pulse, ϕ_2 , to scan through the interference fringe, and obtain the fringe contrasts from the fit. Figure 6.11 summarises the fringe contrasts obtained along both trajectories (red connected dots) as a

²The dynamic range of an interferometer with $2T = 0.2$ ms is around 20 m/s^2 , which exceeds the environmental, vibrational noise of approximately 0.5 mm/s^2 . Hence, the interferometer phase noise due to vibrations is negligible.

function of height. For comparison, the blue dashed line shows the relative beam intensity measured using the π -times of Rabi oscillations. A clear improvement in the fringe contrast is observed away from the centre of the beam. I was able to make further progress in the interferometry by keeping the atoms away from the centre of the Raman beam, but clearly there was a problem with the Raman intensity distribution, which needed eventually to be fixed. In section 6.2.5 I briefly move forward in time to show a very recent result of imaging the Raman beam directly on a camera.

6.2.5 Confirmation with the camera

The unexpected intensity features in the Raman beam have been very recently confirmed by directly taking a series of images of the beam profile, when the chamber was vented. The CCD chip of the camera is quite small, having dimensions 6.5 mm x 4.9 mm but we were able to reconstruct the full beam profile by stitching together over 40 images. Figure 6.12 (a) shows all images, which have been arranged as a mosaic. A bright spot stands out in the middle of the beam, which is located very close to the initial position of the atom cloud at the launch. Figure 6.12 (b) exhibits a closer look of the spot.

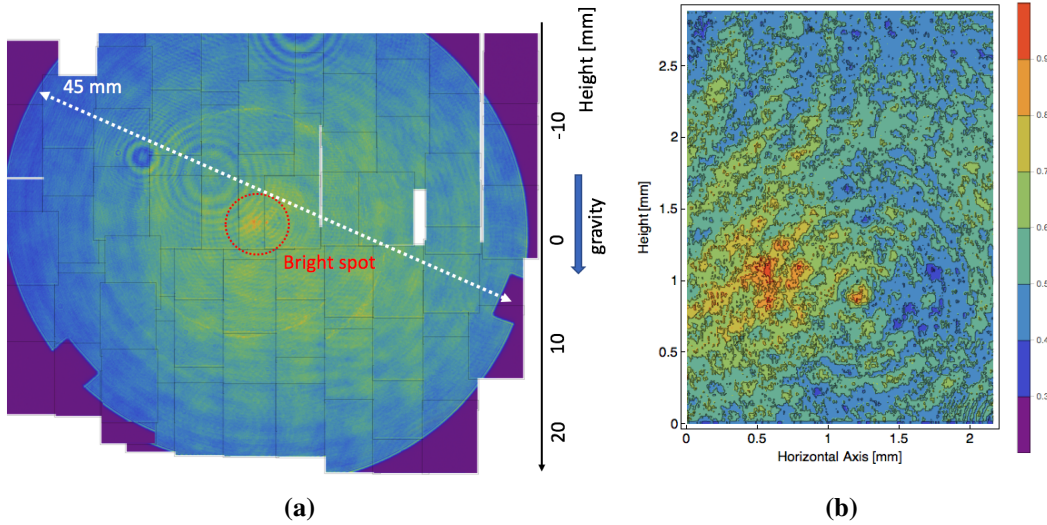


Figure 6.12: Mosaic images of the Raman beam profile measured on a CCD chip (a). A close-up image is shown in (b), where the normalised pixel values vary between 0.23 and 1.

The pixel values in Fig. 6.12 (b) vary between 0.23 to 1, which show intensity features at higher spatial frequency. To resemble the averaged beam intensity seen by the atom cloud, I apply a Gaussian filter with a radius of 2 mm to Fig. 6.12 (b). The smoothed pixel values vary by 20% across the image after filtering. This agrees with the intensity measurements around 0 mm in Fig. 6.7. In addition to the bright spot, it is also worth mentioning that the intensity rings in (a) are the results of the surface impurities of the

purchased aspherical doublet lens pairs, which are made of molten glass. This will also be replaced in the next version of the Raman optics design.

6.3 Contrast vs $2T$ -time

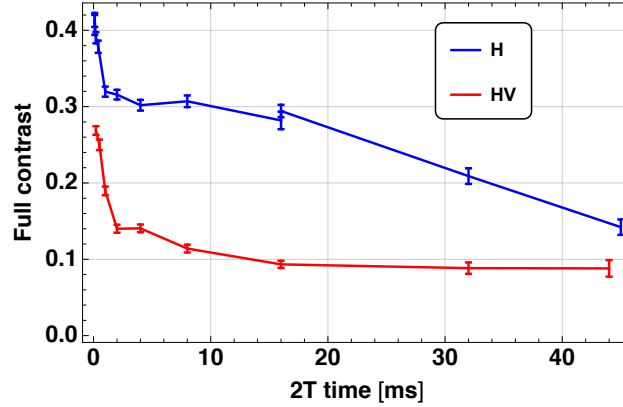


Figure 6.13: Interferometer fringe contrast versus the interferometer time $2T$. Red dots: launching to stay near the centre of the Raman beams. This gives poor contrast because of the defect in the beam profile. Blue dots: launching to avoid the defect region.

After understanding the nature of the problem, as described in sections 6.2.1-6.2.4, we were able to obtain improved fringe contrasts at longer interferometer times T . Figure 6.13 shows the measured contrast as a function of the interferometer time $2T$. The connected red dots show the dramatic loss of contrast with increasing $2T$ when the atoms are launched on the HV trajectory, which keeps them near the centre of the Raman beam, in the vicinity of the defect. This behaviour was already shown in Fig. 6.3. The blue dots, on the other hand, are measured using atoms on the H trajectory, which quickly falls below the centre of the Raman beam. To avoid the defect, these fringes all have the third interferometer pulse operating when the atom are 20 mm below the centre of the Raman beam, where the intensity profile appears to be smoother, as shown in Fig. 6.12 (a). To increase T , we keep the third pulse there, and perform the first two pulses earlier in time. As T increases, the starting height of the interferometer gradually moves upwards the centre of the Raman beam where the intensity is higher. We therefore adjust the first and second interferometer pulse durations for each interferometer, in order to get as close as possible to the desired $\pi/2$ and π pulse areas. Note that there is a very steep decrease in visibility over the first one or two milliseconds, but the subsequent loss of visibility up to 45 ms is relatively slow. The steep decrease is likely due to phase noise in the Raman laser. The slow decrease is a combination of the laser phase noise and the increasing significance of the interferometer phase noise due to environmental vibration. However, further investigation is required to full understand this. Nevertheless, the present chap-

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ter has shown that we can greatly improve the fringe contrast by avoiding the defect, maintaining 30 % of contrast for $2T$ up to 15 ms and 15 % at $2T = 45$ ms.

Chapter 7

Noise

In this chapter I begin by exploring how the acceleration signal derived from the 3-pulse interferometer becomes noisy as a result of noise in the light, the detection, or the electronics, or through real acceleration noise in the environment. Section 7.1 sets up a general theoretical framework for describing the response of the interferometer to various noise sources. Section 7.2 presents experiments to characterise the noise. In section 7.3 I describe experiments to compare the interferometer output with the reading of a classical MEMS accelerometer mounted to the Raman retro-reflection mirror.

7.1 Theoretical treatment of the interferometer response to various noise sources

7.1.1 Response to noise in the optical phase difference

The interferometer phase $\Delta\Phi$ in Eq. 2.23 depends on the phase difference between the two laser beams that drive Raman transitions, which is imprinted on the wave function of the atom. This phase difference is equally affected by the position of the atoms relative to the retro-reflecting mirror and the imperfect laser phase lock. Noise in both results in noise of the interferometer phase, which is equivalent to acceleration noise.

Let us consider a fluctuation $\delta(\Delta\Phi)$ in the interferometer phase, due to small fluctuations over time $d\phi(t)$ in the optical beat note phase between the two laser beams driving the Raman transitions. Following [72, 73], the interferometer phase noise can be written as an integral over time involving $\phi(t)$:

$$\delta(\Delta\Phi) = \int g(t) d\phi(t) = \int g(t) \frac{d\phi(t)}{dt} dt. \quad (7.1)$$

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For a $\pi/2 - \pi - \pi/2$ interferometer with π - and $\pi/2$ -pulse durations being τ_π and $\tau_{\pi/2}$ respectively, the sensitivity function $g(t)$ is [74]

$$g(t) = \begin{cases} -\sin(\Omega_{\text{eff}}t), & 0 < t < \tau_{\pi/2} \\ -1, & \tau_{\pi/2} < t < \tau_{\pi/2} + T \\ -\sin(\Omega_{\text{eff}}(t - T)), & \tau_{\pi/2} + T < t < \tau_{\pi/2} + T + \tau_\pi \\ 1, & \tau_{\pi/2} + T + \tau_\pi < t < \tau_{\pi/2} + 2T + \tau_\pi \\ -\sin(\Omega_{\text{eff}}(t - 2T)), & \tau_{\pi/2} + 2T + \tau_\pi < t < 2\tau_{\pi/2} + 2T + \tau_\pi \\ 0, & \text{elsewhere.} \end{cases} \quad (7.2)$$

Here, T is the time interval when the light is off, between pulse 1 and 2 and between 2 and 3. For short interferometer pulse durations, where $\tau_{\pi/2} \ll T$, we can simplify Eq. 7.2 to be

$$g(t) = \begin{cases} -1, & 0 < t < T \\ 1, & T < t < 2T \\ 0, & \text{elsewhere.} \end{cases} \quad (7.3)$$

Equation 7.1 relates the fluctuation of the interferometer phase to the fluctuations $d\phi(t)$ but it is more convenient to characterise the noise in ϕ by its spectral power density $S_\phi(\omega)$. Reference [75] relates the variance of the interferometer phase $(\sigma_{\Delta\Phi}^{\text{rms}})^2$ to $S_\phi(\omega)$:

$$(\sigma_{\Delta\Phi}^{\text{rms}})^2 = \int_0^{+\infty} |H(\omega)|^2 S_\phi(\omega) d\omega \quad (7.4)$$

via a transfer function $|H(\omega)|^2 = |\omega G(\omega)|^2$, where $G(\omega)$ is the Fourier transform of the sensitivity function in Eq. 7.2. Figure 7.1 shows the transfer function $|H(2\pi f)|^2$ for an interferometer with $2T = 50$ ms with a π -pulse duration of $\tau_\pi = 20$ μs as a function of frequency. At frequencies $f < 1/(2T + 2\tau_\pi)$, the interferometer is insensitive to

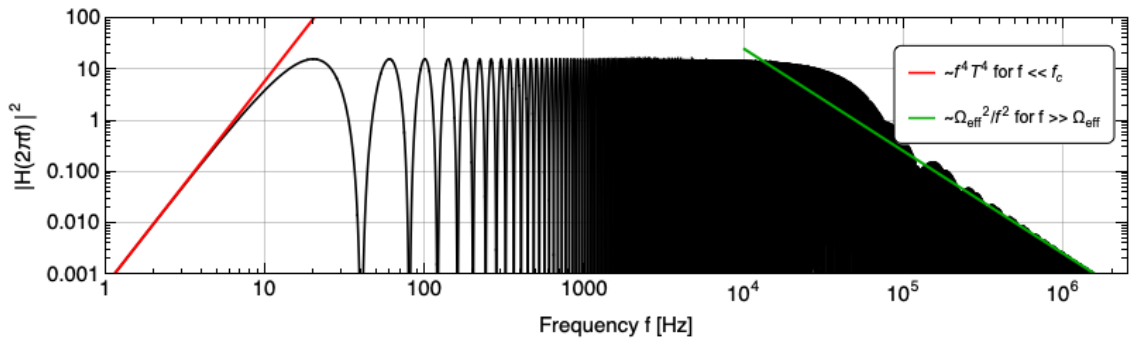


Figure 7.1: Transfer function $|H(2\pi f)|^2$ of a $2T = 50$ ms interferometer with a π -pulse duration of 20 μs .

the slow change of the optical phase difference. In that regime, the transfer function is approximately proportional to $f^4 T^4$, as illustrated in the red line in Fig. 7.1. As the frequency increases, the transfer function oscillates periodically, with zeros at integer multiples of $1/(T + \tau_\pi)$, shown in the black solid line. When the frequency exceeds the Rabi frequency Ω_{eff} , the optical phase fluctuations average out during each laser pulse, and then the transfer function behaves as a low-pass filter, indicated by the green line in Fig. 7.1.

As the interferometer performs a sequence of discrete acceleration measurements at a cycling rate of f_c , the Allan variance is a useful tool to characterise the stability of the interferometer over N successive measurements. The time taken for N measurements is simply N/f_c and the corresponding mean value of the interferometer phase over the N measurements is denoted by $\langle \Delta \Phi_N \rangle$. Following Allan [76], we define the Allan variance of the interferometer phase $\sigma_{\Delta \Phi}^2$ as a function of N by

$$\sigma_{\Delta \Phi}^2(N) = \frac{1}{2} \lim_{\mathcal{N} \rightarrow \infty} \left(\frac{N}{\mathcal{N}} \sum_{k=1}^{\mathcal{N}/N} (\langle \Delta \Phi_N \rangle_{k+1} - \langle \Delta \Phi_N \rangle_k)^2 \right), \quad (7.5)$$

where $\mathcal{N}$ indicates the total number of acceleration measurements acquired. Equation 7.5 describes the average variance between adjacent mean interferometer phases, where the means are taken over groups of N measurements with the group index k .

For large N , Ref. [74, 75, 77] show that the Allan variance in Eq. 7.5 can be related to the transfer function of the interferometer and the power spectral density of the phase of the optical beat note, and is given by

$$\sigma_{\Delta \Phi}^2(N) = \frac{f_c}{N} \sum_{n=1}^{\infty} |H(2\pi n f_c)|^2 S_\phi(2\pi n f_c). \quad (7.6)$$

Equation 7.6 shows that the interferometer is only sensitive to the optical beat note phase noise at integer multiples of the cycling frequency, $n f_c$. As a time gap exists between subsequent interferometers, which is used for MOT loading, cooling and state preparation, this leads to a degradation of the sensitivity of the interferometer due to aliasing effect, which was first recognised as the Dick effect [72, 73] in atomic clocks. Section 7.2.3 looks at the increase of interferometer's acceleration noise in the presence of dead time.

7.1.2 Response to ambient vibrations

In addition to the noise arising from an imperfect phase lock, there can be jitter in the optical phase difference seen by the atoms, resulting from the vibration of the retro-reflecting

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mirror. We can replace the optical phase fluctuation in Eq. 7.1 by $d\phi = 2k_2 dx(t)$, where $dx(t)$ is the fluctuation in the distance between the atom and the mirror. Using

$$\frac{d\phi(t)}{dt} = 2k_2 \frac{dx(t)}{dt} = 2k_2 v(t), \quad (7.7)$$

where $v(t)$ is the velocity of the mirror, the interferometer phase fluctuation accumulated over one interferometer cycle can be rewritten as

$$\delta(\Delta\Phi) = 2k_2 \int_0^{2T+2\tau_\pi} g(t) v(t) dt. \quad (7.8)$$

For a single measurement, Eq. 7.2 shows that $g(t)$ vanishes for $t \notin [0, 2T + 2\tau_\pi]$. Hence, we can extend the boundaries of the integral in Eq. 7.8 to $\pm\infty$, which can then be solved using integration by parts [74]:

$$\delta(\Delta\Phi) = -2k_2 [f(t) v(t)]_{-\infty}^{+\infty} + 2k_2 \int_{-\infty}^{+\infty} f(t) a(t) dt, \quad (7.9)$$

where $a(t)$ is the acceleration of the distance between the atoms and the mirror and

$$f(t) = - \int g(t') dt'. \quad (7.10)$$

For our interferometer with $\tau_\pi \ll T$, $f(t)$ can be approximated by [74]

$$f(t) = \begin{cases} t, & 0 < t < T \\ T - t, & T < t < 2T \\ 0, & \text{elsewhere.} \end{cases} \quad (7.11)$$

As $f(t)$ also vanishes for $t \notin [0, 2T]$, the first term in Eq. 7.9 can be ignored and the interferometer phase fluctuation is reduced to

$$\delta(\Delta\Phi) = 2k_2 \int_{-\infty}^{+\infty} f(t) a(t) dt. \quad (7.12)$$

The constant part of $a(t)$ is, of course, the acceleration that we aim to measure, so it is the fluctuating part of $a(t)$ that generates the noise. Analogously to the last section 7.1.1, the variance of the interferometer phase can be related to the power spectral density of the acceleration, denoted S_{acc} , and the Fourier transform of the response function $F(\omega)$:

$$(\sigma_{\Delta\Phi}^{\text{rms}})^2 = 4k_2^2 \int_0^{+\infty} |F(\omega)|^2 S_{\text{acc}}(\omega) d\omega. \quad (7.13)$$

7.1. THEORETICAL TREATMENT OF THE INTERFEROMETER RESPONSE TO VARIOUS NOISE SOURCES

The transfer function $|H_{\text{acc}}(\omega)|^2$ connecting the interferometer variance to S_{acc} is closely related to the $|H(\omega)|^2$ in Eq 7.4, as proved in Ref. [75]:

$$|H_{\text{acc}}(\omega)|^2 = 4k_2^2 |F(\omega)|^2 = 4k_2^2 \frac{|H(\omega)|^2}{\omega^4}. \quad (7.14)$$

Figure 7.2 shows this function for a $2T = 50$ ms interferometer. For frequencies below $(2T + 2\tau)^{-1}$, the transfer function is almost flat, while it rolls off at high frequencies as ω^{-4} .

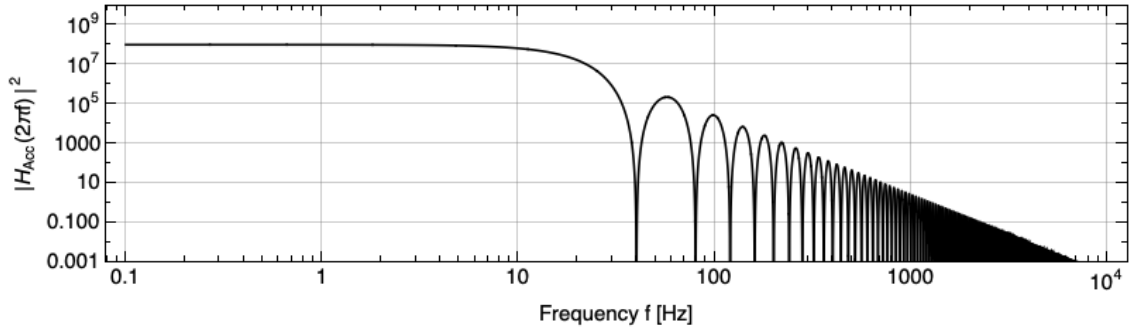


Figure 7.2: Transfer function to acceleration for an interferometer with $2T = 50$ ms.

If the power spectral density of the acceleration noise S_{acc} is known, the Allan variance of the interferometer phase due to mirror vibration is described for large N , analogously to Eq. 7.6, as

$$\sigma_{\Delta\Phi}^2(N) = \frac{f_c}{N} \sum_{n=1}^{\infty} |H_{\text{acc}}(2\pi n f_c)|^2 S_{\text{acc}}(2\pi n f_c), \quad (7.15)$$

These results will be used in section 7.2 to study the noise in our interferometer.

7.1.3 Quantum projection noise

Fluctuations of the beat note phase are not the only source of interferometer noise. There is also noise inherent in the readout of the interferometer. As described in section 3.8.1 after the third interferometer pulse, the probability $P_{F=2} = N_2/N_0$ is read out by measuring the number of $F = 2$ atoms (N_2) and then the total number of atoms (N_0). Quantum projection noise [78] in the value N_2 introduces uncertainty into this measurement, which gives $P_{F=2}$ a standard deviation of

$$\sigma_{\text{QPN}} = \sqrt{\frac{P_{F=2}}{N_0} (1 - P_{F=2})}. \quad (7.16)$$

The interferometer is most sensitive to acceleration at the points of greatest slope, where $P_{F=2} = 1/2$. Assuming it is operating there, the quantum projection noise in the measured $P_{F=2}$ is $\sqrt{\frac{1}{4N_0}}$. Equation 2.29 shows how to relate that to an equivalent standard deviation of the acceleration. In Chapter 6 we achieved $C = 0.15$ with $T = 25$ ms and 10^6 atoms, so the quantum projection noise corresponds to a standard deviation in acceleration of $0.7 \mu\text{m/s}^2$. This is comfortably better than our design of 100 ng uncertainty at 1 second, so it is not a limiting problem.

It is worth mentioning that there are some atoms remaining in $|1, \pm 1\rangle$ states after the state preparation described in section 3.8.2. These contribute to the detection of the total number of atoms during the state sensitive detection (see section 3.8.1) to estimate the $F = 2$ probability. As a correction, we measure the atom number in $|1, \pm 1\rangle$ in an auxiliary measurement and subtract it from N_0 . Appendix 9.2 summarises the contribution to the interferometer phase noise due to the noise in the number of atoms in $|1, \pm 1\rangle$, which is not significant as that number is stable.

7.1.4 Shot noise

The number of atoms in the MOT fluctuates from shot to shot, which is canceled by taking the ratio between N_2 and N_0 . However, both number measurement involve the counting of photons, which has Poissonian statistics. With n photons detected per atom, the standard deviations are $\sqrt{nN_2}$ and $\sqrt{nN_0}$. We use error propagation to write $\sigma_{P_{F=2}}$ due to photon shot noise as

$$\sigma_{\text{ph}} = \sqrt{\frac{P_{F=2}}{nN_0} (1 + P_{F=2})}. \quad (7.17)$$

In my atom interferometer, because $n \gg 1$, the number of detected photons will be always larger than the atom number, and therefore it is safe to take the quantum projection noise to be always larger than the photon shot noise, $\sigma_{\text{QPN}} > \sigma_{\text{ph}}$.

7.2 Experimental study of noise

7.2.1 Detector electronics noise

For the interferometer to operate with the lowest possible noise level, we want it to be limited by the quantum projection noise. For $N_0 = 10^6$ atoms in the interferometer with $P_{F=2} = 1/2$, the number of atoms in $F = 2$ is given by $N_2 = P_{F=2}N_0$. When illu-

minated by two opposing imaging beams with a detuning of -3 MHz^1 at the saturation intensity ($s = 1$), the photodiode collects around 8 nW of optical power from the $F = 2$ atoms. The quantum projection noise causes a fluctuation in the number of atoms N_2 with a standard deviation of $\sigma_{\text{QPN}} N_0$ where σ_{QPN} is given in Eq. 7.16. The noise in the optical power due to the quantum projection noise is around 8 pW . In the apparatus, we

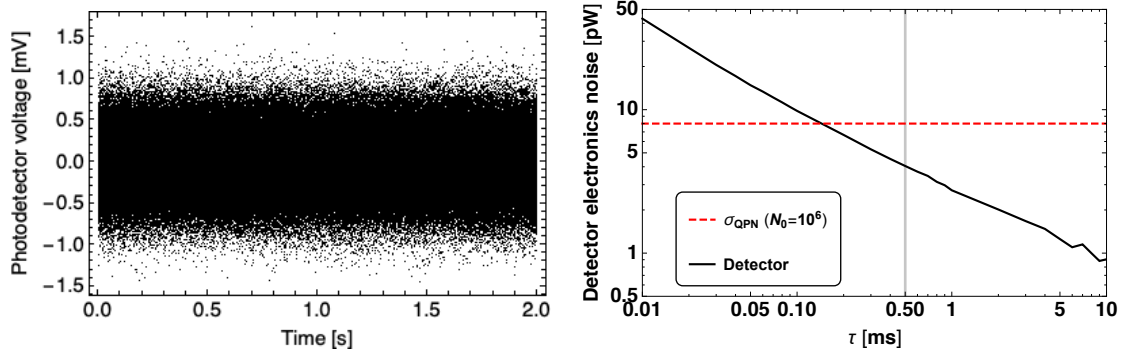


Figure 7.3: Output voltage of the Femto photodiode (black) (a) without any light incident on it. (b) Allan deviation of the detector electronics noise in pW (black) as a function of the duration when the photodiode is on. The dashed red line shows the quantum projection noise for 10^6 atoms.

use a low noise photodiode (*FEMTO* LCA-S-400K-SI). I recorded the output voltage of this photodetector as a function of time without any light being incident on it in order to measure the instrumental noise. From a trace shown in Fig. 7.3 (a) taken over 2 seconds with a resolution of $5 \mu\text{s}$, I computed the Allan deviation plotted by the black line in Fig. 7.3 (b) as a function of the averaging time τ . For comparison, the red dashed line indicates the noise in the optical power due to the quantum projection noise for $N_0 = 10^6$ atoms. The integration time that I have been using is $\tau = 500 \mu\text{s}$, shown by the vertical grey gridline. This puts the detector noise about a factor of 2 below the quantum projection noise, which is a satisfactory noise level. It could be lowered further by increasing the integration time τ , but then atoms start to be optically pumped out of $F = 2$ because the detection light is not perfectly polarised, and this begins to compromise the accuracy of determining the interferometer phase.

Because the detector noise comes into the measurement of N_0 as well as N_2 , let us now consider the contribution of the detector noise to the interferometer phase noise $\sigma_{\Delta\Phi}$. For simplicity, we assume all other noise sources are negligible. Analogously to Eq. 2.28, the interferometer phase noise due to the detector noise $\sigma_{\Delta\Phi}^{(D)}$ is described by:

$$\sigma_{\Delta\Phi}^{(D)} = \sigma_{P_{F=2}}^D / \frac{dP_2}{d\Phi} = \frac{2}{C} \sigma_{P_{F=2}}^D, \quad (7.18)$$

¹To prevent atoms escaping the detection area due to heating during the imaging, the imaging frequency is detuned by -3 MHz from the $|F = 2\rangle \leftrightarrow |F' = 3\rangle$ transition.

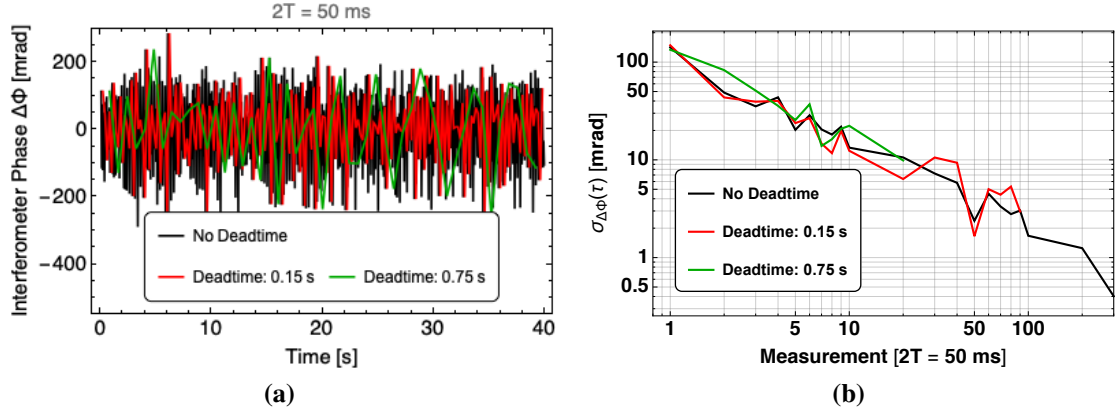


Figure 7.4: The interferometer phase $\Delta\Phi$ of a $2T = 50$ ms interferometer due to variations in the Raman laser phase difference is plotted in (a) as a function of time. The interferometer phases are calculated using the data shown in Fig. 4.2 (a). Three different colours show the interferometer phases estimated without any dead time (black), with 0.15s (red) and 0.75 s (green) of dead time between the interferometers. (b) shows the Allan deviations of the interferometer phase in (a) as a function of the number of measurements.

where $\sigma_{P_{F=2}}^D$ is the noise of $P_{F=2}$ only due to the detector noise. Using $P_{F=2} = V_2/V_0$, where V_2 and V_0 are the photodiode voltage obtained when detecting N_2 and N_0 , we describe $\sigma_{P_{F=2}}^D$ as:

$$\sigma_{P_{F=2}}^D = \sqrt{\frac{1}{V_0^2}\sigma_V^2 + \frac{V_2^2}{V_0^4}\sigma_V^2} = \frac{\sigma_V}{V_0} \sqrt{1 + P_{F=2}^2}, \quad (7.19)$$

with the noise of the detector being σ_V . In summary, Eq. 7.18 turns into

$$\sigma_{\Delta\Phi}^{(D)} = \frac{2}{C} \frac{\sigma_V}{V_0} \sqrt{1 + P_{F=2}^2}. \quad (7.20)$$

For 10^6 atoms with a photodiode voltage of $V_0 = 63$ mV and a detector noise of $\sigma_V = 0.02$ mV, $\sigma_{\Delta\Phi}^{(D)}$ is expected to be around 5 mrad in the middle of the fringe, where the contrast is assumed to be $C = 0.15$. This corresponds to an acceleration noise of $0.5 \mu\text{m/s}^2$ for an interferometer operating at $2T = 50$ ms. In contrast, the interferometer phase noise due to the quantum projection noise is around 7 mrad, which is equal to an acceleration noise $0.7 \mu\text{m/s}^2$.

7.2.2 Laser phase noise

Figure 4.2 in section 4.1 shows a measurement of variation in the phase of the beat note between the two Raman frequencies. Here I use this to estimate the resulting contribution to the interferometer phase noise. Firstly, the data in Fig. 4.2 (a) (top) are partitioned into datasets of a length of 50 ms. For each set I calculate the interferometer phase $\Delta\Phi$ of a $2T = 50$ ms interferometer using Eq. 7.1. Taking only every n -th dataset corresponds to

operating the interferometer with a dead time of $(n - 1) \times 50$ ms. Figure 7.4 (a) shows the calculated interferometer phases by taking every datasets (black), every 4th (red) and every 16th (green), which are equivalent to having dead times of 0 ms, 150 ms and 750 ms respectively. Figure 7.4 (b) summarises the corresponding Allan deviations of the calculated interferometer phase noise $\sigma_{\Delta\Phi}$, shown in Fig. 7.4 (a), as a function of the number of measurements. It is interesting to see that the contribution of the phase noise of the beat note between the Raman lasers to the interferometer phase is not affected by the dead time between measurements. The same levels of interferometer phase noise are found after performing the same number of measurements, regardless of the time gap between the measurements.

The Allan deviations in (b) show a clear $1/\tau$ dependence. This dependence is typical of the noise produced by the rounding errors in a digitised signal. Further investigation of the data acquisition of the beat note phase between the two Raman frequencies is needed to determine if that is the case here. I therefore take the noise presented in Fig. 7.4 as an upper bound on the laser phase noise.

7.2.3 Vibration noise

The acceleration noise due to ambient vibration can be bigger than the fringe spacing Δa . A pendulum MEMS accelerometer fixed to the back of the retro-reflecting mirror (see section 4.3) allows us to know the order number of the atom interferometer fringe. The apparatus sits on a laser table that can be floated on pneumatic legs for vibration isolation. Figure 7.5 shows time series of the acceleration measured by the pendulum accelerometer when the table is floated (red) and not floated (blue). The data in Fig. 7.5 (a) are sampled at 100 kHz, giving standard deviations of the acceleration with the floating and non-floating table of approximately 1 mm/s^2 and 3 mm/s^2 over 40 s of measurement. Figure 7.5 (b) exhibits the acceleration measured over 10^4 s, where the data points are each averaged over one second. Both traces in Figure 7.5 (b) show similar drifts of acceleration of approximately $100 \text{ nm/s}^2/\text{s}$ over the course of the measurement. We do not know the origin of this drift, but surmise that it may simply be the drift of the MEMS accelerometer zero (bias). In Fig. 7.6, the red and blue curves represent the power spectral density (PSD) of the acceleration noise shown in Fig. 7.5 (b), when the table is floated and not floated. We see that the acceleration noise is heavily suppressed between 5 and 200 Hz when the pneumatic isolation system is switched on. The peaks around 2 Hz in the red line represent the resonant frequencies of the pneumatic isolation system. In comparison, the black curve shows the (much lower) noise spectrum of the *National Instruments* data acquisition card (NI card - PXI 4462) used to read the accelerometer. The effect of the ambient vibrations is also visualised in the Allan deviations. The solid lines in

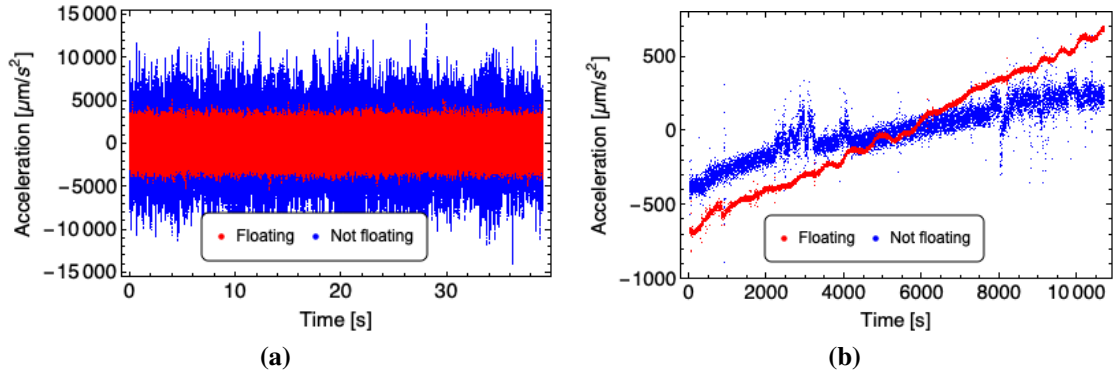


Figure 7.5: Ambient vibration measurement using the pendulum MEMS accelerometer in vacuum. (a) Time series of acceleration measured over 40 s at a sampling rate of 100 kHz. (b) Time series taken over 10000 s and averaged every second.

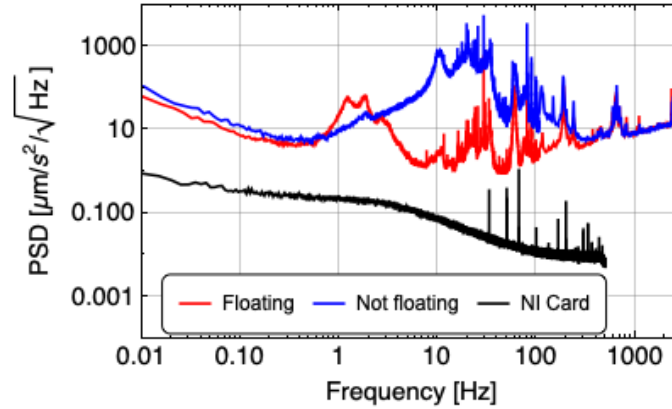


Figure 7.6: Vibrational power density spectra measured by the classical pendulum accelerometer with the pneumatic isolation platform being switched on (red) and off (blue). The black curve illustrates the noise floor of the acquisition card.

Fig. 7.7 exhibit the Allan deviations of the acceleration data presented in Fig. 7.5. As the bandwidth of the MEMS accelerometer is only 2 kHz, acceleration noise at $\tau < 0.25$ ms is a result of the noise in the electronics in the accelerometer. That explains why the Allan deviations are the same in that range regardless of whether the table is floating or not floating. As τ increases, the MEMS accelerometer becomes sensitive to the real vibrations. The deviation reduces by a factor of ten when the table is floated. The lump in the solid blue line around $\tau = 10$ ms is due to the noise between 10 to 200 Hz in Fig. 7.6, whereas the lump in the solid red line around 0.3 s is caused by the resonance frequency of the floating table at around 2 Hz. At a correlation time of about 5 seconds, the floating table has a minimum noise of $4 \mu\text{m/s}^2$. Both curves increase with time beyond 10 s, with the blue curve exhibiting a short range of approximately $\tau^{1/2}$ dependence, which could indicate some random walk [79] of the acceleration. For $\tau > 1000$ s, both curves show a τ^1 -dependence. This is characteristic of a drift in the acceleration. The most likely culprit

is the bias drift, which is sensitive to temperature at a rate of $0.3 \text{ mm/s}^2/^\circ\text{C}$. The drift we see here is approximately $0.3 \text{ mm/s}^2/\text{hour}$, which would indicate a temperature change of approximately 1 deg/hr . That is entirely consistent with the temperature instability in our laboratory.

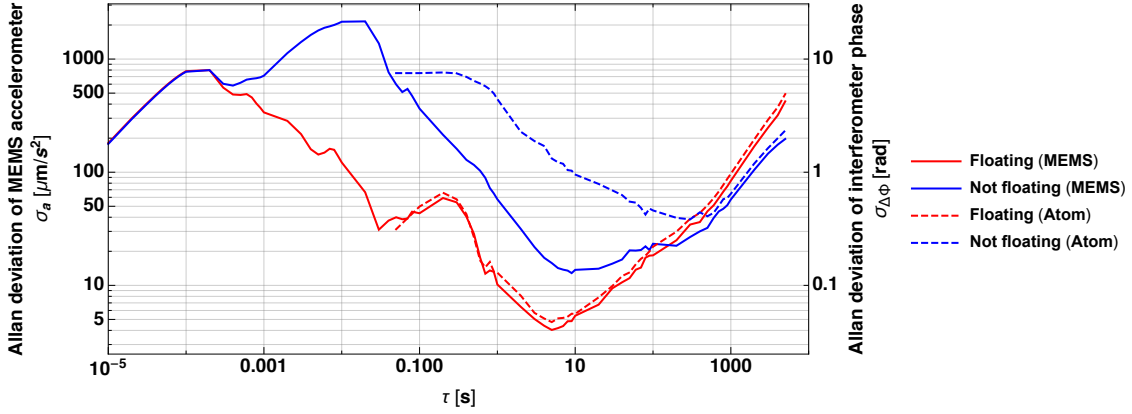


Figure 7.7: Allan deviation of the acceleration measured by the MEMS accelerometer when the table is floating (solid red) and not floating (solid blue). The dashed lines show the deviations of acceleration noise in a 50 ms interferometer calculated from the data presented in Fig. 7.5.

Because the atom interferometer does not sample the acceleration continuously, it is an interesting exercise to determine how the phase of the atom interferometer would respond to this level of noise. (For this purpose, we do not care whether the slow drift is real or not.) Equation 7.12 transforms the time-dependent acceleration into the time-dependent interferometer phase, from which I calculate the Allan deviations shown as dashed lines in Fig. 7.7. I have taken $2T = 50 \text{ ms}$, and the scale is on the right. For the case of the floating table, the Allan deviation of the atom interferometer phase shown in the dashed red line perfectly follows the underlying noise level (solid red line). This shows that the atom interferometer accurately captures the level of noise of the floating table. In contrast for the case when the table is not floating, the Allan deviation of the interferometer noise (dashed blue line) is significantly higher than the level of underlying noise (solid blue line) between $\tau = 0.1 \text{ s}$ and 100 s . This is due to aliasing of the vibrations in the environment between 10 and 100 Hz , which are shown by the blue peak in Fig. 7.6. This extra noise has also caused the correlation time with the lowest acceleration noise in the blue lines to be extended from 10 s to few hundreds seconds. At longer correlation time τ , the phase noise of the atom interferometer is again in very good agreement with the expectation, as the noise at shorter correlation time are all averaged out. It is worth mentioning that aliasing is less obvious for the case of the floating table, as the vibrational noise at higher frequency is by two orders of magnitude smaller than with the table not-floating, and hence the aliased noise is negligible.

The interferometer Allan deviation in Fig. 7.7 is calculated from the MEMS signal

7.3. COMPARISON BETWEEN THE ATOM INTERFEROMETER AND THE CLASSICAL ACCELEROMETER

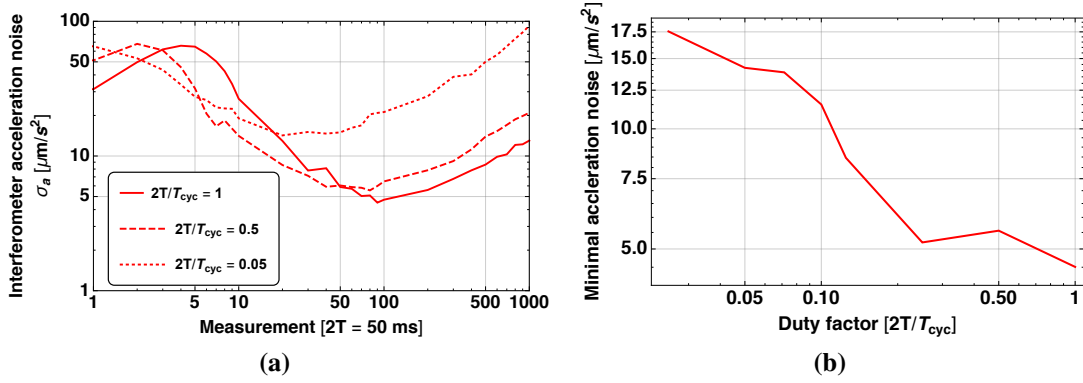


Figure 7.8: Allan deviation of the MEMS accelerometer readings as a function of the number of measurements for interferometers operating at $2T = 50$ ms (a). The solid line shows the Allan deviation in the absence of dead time, whereas the dashed and dotted lines indicate the acceleration noise for duty cycles of 0.5 and 0.05. (b) shows an increase of the minimal acceleration noise as a function of the duty factor.

without any dead time between measurements. However, in real life, the time gaps do exist, as they are required for MOT loading, cooling and state preparation. In the following, we saw above that this is not important for the phase noise in the laser beat note, but here I consider briefly the effect of dead time on the noise arising from the MEMS accelerometer reading. Figure 7.8 (a) shows the averaged interferometer acceleration noise (with $2T = 50$ ms) as a function of the number of measurements being averaged, as calculated from the MEMS accelerometer readings on the floating platform. The solid red line in Fig. 7.8 (a) is the same as the dashed red line in Fig. 7.7, where there is no dead time. In this case, the cycling time T_{cyc} between measurement sequences is equal to the interferometer time $2T$. The dashed and dotted lines in Fig. 7.8 (a) show the same calculations with duty factor $2T/T_{\text{cyc}}$ of 0.5 and 0.05 respectively. As we increase the dead time, the same number of measurement corresponds to a longer time, and this makes the curves in Fig. 7.7 move to the left. In addition, the linear drift accumulates a larger error between readings, with the result that the minimum of the curve increases. Figure 7.8 (b) plots that minimum as a function of the duty factor. Even for a cycling time of 1 s and hence a duty factor of 0.05 at $2T = 50$ ms, the noise minimum is only $15 \mu\text{m/s}^2$. This is quite small enough to identify the order number of the interferometer fringe reliably, given that the fringe spacing is $\lambda/(2T^2) = 624 \mu\text{m/s}^2$.

7.3 Comparison between the atom interferometer and the classical accelerometer

In section 7.1.2, we theorised about how the atom interferometer would respond to the acceleration of the apparatus, as measured by the MEMS accelerometer reading $a(t)$.

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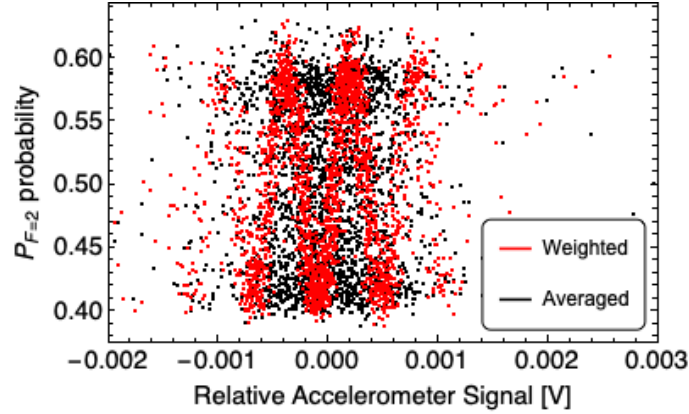


Figure 7.9: Measurements of $P_{F=2}$ plotted against the MEMS accelerometer. Black points: MEMS voltage is averaged over 45 ms of the interferometer measurement. Red points: MEMS voltage is averaged using Eq. 7.12 with the triangular weighting function in Eq. 7.11.

Equation 7.12 was used to derive the expected interferometer phase $\Delta\Phi$ from $a(t)$, in which the transfer function $f(t)$ describes how the interferometer responds to a changing acceleration. We saw in Eq. 7.11 that $f(t)$ is well approximated as a simple triangular weighting function. In this section, I describe an experiment to test the usefulness of this weighting by comparing the measured output of the atom interferometer with the triangularly-weighted $a(t)$ and with a simple top-hat averaged $a(t)$. With a fixed phase of the Raman beat note, we measure over 40 minutes both the $P_{F=2}$ output of an interferometer operating at $2T = 45 \text{ ms}^2$ and the signal of the pendulum accelerometer, which is sampled at 200 kHz. Figure 7.9 shows $P_{F=2}$ as a function of the pendulum accelerometer signal. For the black points, each interferometer reading is plotted against the unweighted average of $a(t)$, taken over the course of the 45 ms interferometer reading. Although the table is vibrating - indeed it is not floating - and although one can see a clustering of the points at the turning points of a fringe pattern, no fringe can be discerned. That is because the simple average of $a(t)$ is not what the interferometer senses. By contrast, when each $P_{F=2}$ is plotted against the triangularly weighted average (red points) the fringe pattern is immediately revealed. The difference is less dramatic when the table is floating because the lower noise level produces a smaller difference between the mean of $a(t)$ and the triangularly weighted value. Still, it is now our normal practice to use the triangular weighting.

²To avoid the intensity lumps at the centre of the Raman beam profile, the interferometer is operated at a later flight time. This gives us no more than 45 ms before the atoms leave the Raman beam.

7.3.1 Synchronising the MEMS accelerometer with the interferometer

The light that drives the Raman transitions responds (essentially) instantly to the vibrations of the mirror, whereas the MEMS accelerometer has a non-zero response time. By delaying the start of $f(t)$ relative to the first interferometer pulse, a better synchronisation is achieved between the two measurements. We can see this as a reduction in the noise on the interference fringe pattern. Figure 7.10 (a) once again shows the measured $P_{F=2}$

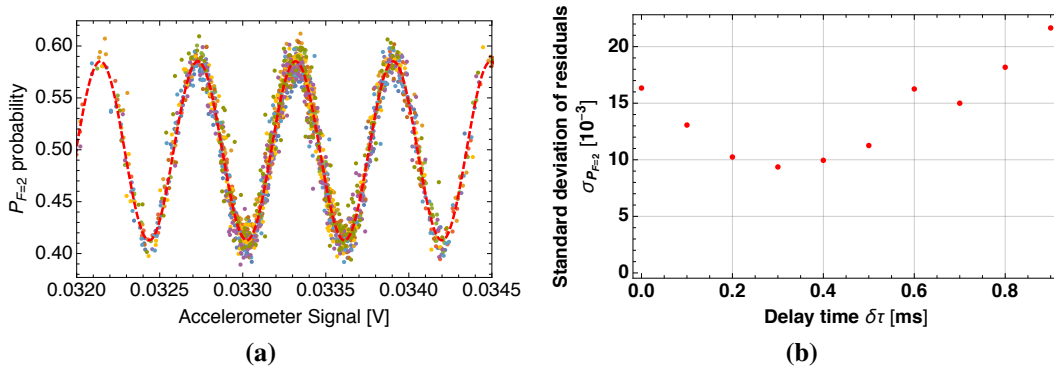


Figure 7.10: (a) Coloured dots show the data points of interferometer fringes with the accelerometer signals collected at different delay times. Solid red line is the fit on all the data points. (b) Residual noise $\sigma_{P_{F=2}}$ plotted versus the delay time for the accelerometer signal, appears to be 0.3 ms at the optimum.

plotted against the triangularly weighted pendulum accelerometer signals with the table not floating. It exhibits 10 independent data sets, where each set is collected for 2 minutes with the MEMS accelerometer data delayed by $\delta\tau$. The colours of the dots indicate the choice of the delay time between 0 to 0.9 ms. The dashed red line is a sinusoidal fit to all the data using

$$P_{F=2} = \frac{1}{2}(P_0 + C \cos(\alpha(V + V_0))). \quad (7.21)$$

The calibration factor α of the pendulum accelerometer will be described in section 7.3.4 in more detail. For each delay time, we calculate the standard deviation $\sigma_{P_{F=2}}$ of the residuals from this fit, and these are plotted against the accelerometer acquisition delay time, $\delta\tau$, in Fig. 7.10 (b). A clear minimum of the noise is shown at a delay time of 0.3 ms, which gives almost a factor of 2 of improvement over having no delay.

7.3.2 Effect of lab lighting

A laser line filter at 780 nm is mounted in front of the photodiode to reduce the room light falling onto the photodiode, when collecting the fluorescence from atoms after the interferometry. To check whether the room light doesn't affect the atom interferometer

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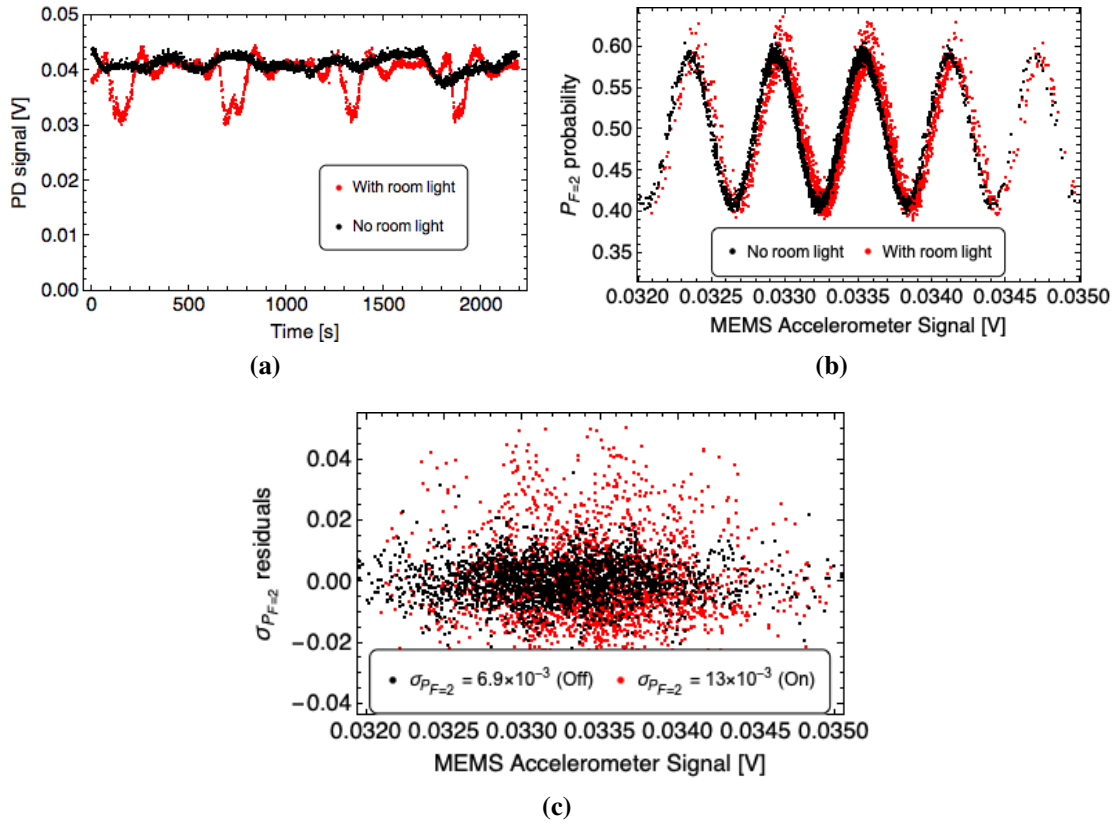


Figure 7.11: Photodiode signal showing V_{tot} measurement over time with the room light on and off respectively (a). The periodic spikes occur on average around every 9 minutes, which vanishes with the room lights off. (b) shows the interference fringe scanned by vibrations, where a clear reduction of noise could be seen in the residual plot (c).

signal, we operated the interferometer with the lab lights switched on and off. Figure 7.11 (a) shows the fluorescence from all the atoms after every interferometer, as recorded by the photodiode over a period of 2200 seconds. The red signal has a recurrent feature with a period of approximately 9 minutes, which vanishes when the room lights are switched off (black). Figure 7.11 (b) exhibits the interference fringes with the room light switched on and off, by plotting $P_{F=2}$ as a function of the MEMS accelerometer signal. One can see a clear difference in the noise level, which is even more evident in the residuals plotted in Fig. 7.11 (c). In addition, the output signal of the pendulum accelerometer is offset by 0.3 mV. If this offset voltage is due to a change in acceleration, we expect an additional 3.3 rad of the fringe to be filled by data points. In contrast, we have measured a phase shift between the fringes to be ≈ 0.8 rad. We are very surprised by these discoveries. The origins of these are still unclear to us. By simply switching off the room light, the noise in the residuals reduces by a factor of 2. Hence, we only operate the interferometer without the room lights. The next student will need to return to this question.

7.3.3 Fringe noise distribution

Each source of noise produces a characteristic variation of noise across the fringe pattern in $P_{F=2}$. Specifically, quantum projection makes a noise $\sigma_{\text{QPN}} \propto \sqrt{P_{F=2}(1 - P_{F=2})}$ (see Eq. 7.16), whereas the detector electronics noise makes $\sigma_D \propto \sqrt{1 + P_{F=2}^2}$ (see Eq. 7.20) and noise in the interferometer phase makes the noise $\sigma_P(\Delta\Phi) \propto \sin(\Delta\Phi)$ (see Eq. 2.28). Note that the noise in the Raman beat note phase and the noise in the MEMS accelerometer both contribute to $\sigma_P(\Delta\Phi)$. These two noise sources are indistinguishable³ in our atom interferometer.

To observe these characteristic noise across the fringe pattern, we operate the interferometer with $2T = 45$ ms at a constant Raman beat note phase and use the vibrations in the environment to scan the fringe pattern. As seen earlier, the vibrational noise sensed by the atoms can be dramatically reduced by floating the table. Figure 7.12 (a) shows two histograms of the horizontal accelerations measured by the classical accelerometer in both environments. The gray histogram was measured with the laser table sitting hard on its legs (grounded), while the green one shows the data with the table floating. Gaussian fits to these histograms give a standard deviation of $637 \mu\text{m/s}^2$ ($\sigma = 0.48$ mV), when the table is grounded and $53 \mu\text{m/s}^2$ ($\sigma = 0.04$ mV), when the table floats. Figure 7.12 (b) shows the corresponding data from the atom interferometer operating at $2T = 45$ ms, plotted as a function of the pendulum accelerometer signal, for which the fringe period is $\lambda/(2T^2) = 770 \mu\text{m/s}^2$. When the table is not floating, the larger vibrational noise moves the point of operation over more than one fringe, as shown in grey dots in Fig. 7.12 (b). The green data points show the same experiment but with the table floating, and there we see that the noise is low enough to keep the operating point near the sensitive part of the fringe. The dashed blue and orange lines are the fits to the data using Eq. 7.21. To observe the variations of the noise across the fringe pattern, let us look at the grey data points.

In Fig. 7.12 (c), the residuals of the fit (dashed blue line in (b)) are also displayed in grey as a function of the same accelerometer signal. The dashed vertical gridlines indicate the positions of greatest slope in the fringes, and there the residuals are clearly larger than elsewhere, as expected for noise in the interferometer phase. Note that noise in the actual acceleration is taken out by plotting $P_{F=2}$ against the MEMS accelerometer reading - indeed, the variation of the real acceleration is the reason that the plot spans several fringes. To estimate the local standard deviations, the fit residuals of $P_{F=2}$ are partitioned into bins with a size of 0.01 mV in the accelerometer voltage. For each bin, I calculate the standard deviation of the $P_{F=2}$ residuals, as shown by the red dots in (c).

³One way to measure the noise in the Raman beat note phase is to use co-propagating Raman beams, so that the interferometer becomes insensitive to acceleration.

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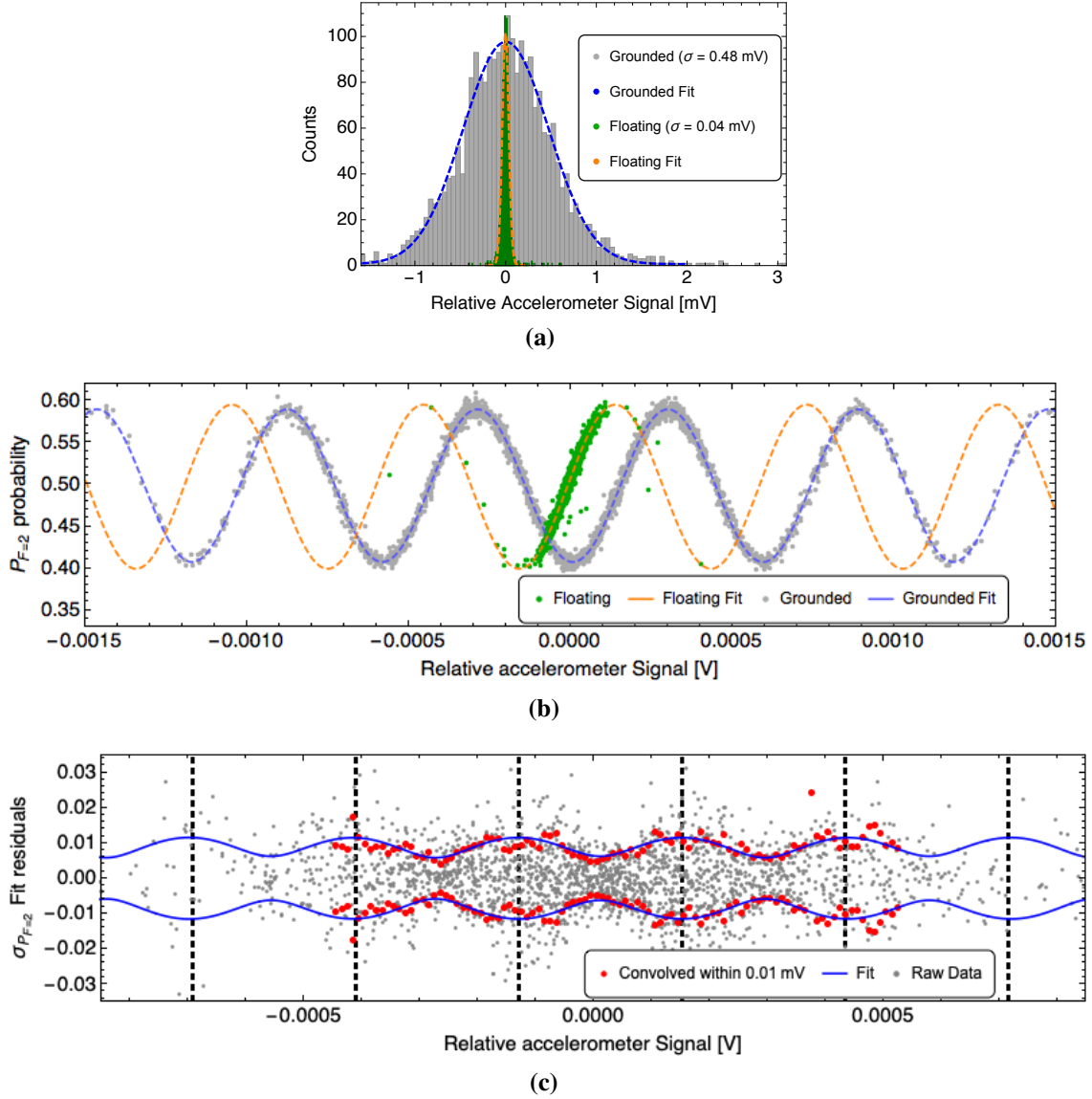


Figure 7.12: Histograms in (a) show the acceleration noise measured by the pendulum accelerometer with the table grounded (gray histogram) and floating (green histogram). (b) shows the fringe patterns measured using an interferometer with $2T = 45$ ms with a $80 \mu\text{s}$ velocity selection pulse in these environments. (c) shows the fit residuals (grey) and a moving averaged residuals (red) with the table not floating. The solid blue line exhibits the fit of the red points using Eq. 7.22.

Assuming all noise sources are uncorrelated, we can add them in quadrature to give

$$\begin{aligned}
 \sigma_{P_{F=2}} &= \sqrt{\sigma_{\text{QPN}}^2 + \sigma_D^2 + \sigma_P^2(\Delta\Phi)} \\
 &= \sqrt{\frac{P_{F=2}}{N_{\text{tot}}} (1 - P_{F=2}) + \frac{\sigma_V^2}{V_{\text{tot}}^2} (1 + P_{F=2}^2) + \frac{C^2 \sin^2(\Delta\Phi)}{4} \sigma_{\Delta\Phi}^2}.
 \end{aligned} \tag{7.22}$$

Although the fringe pattern is plotted as a function of the MEMS accelerometer voltage, we convert that to interferometer phase $\Delta\Phi$ using the fit shown in Fig. 7.12 (b). That al-

7.3. COMPARISON BETWEEN THE ATOM INTERFEROMETER AND THE CLASSICAL ACCELEROMETER

allows me to fit the red dots in Fig. 7.12 (c) to Eq. 7.22 in order to determine the magnitudes of the various noise sources. At the point of greatest slope, the projection noise is only $\sigma_{\text{QPN}} = 0.0006$, which corresponds to an equivalent interferometer phase noise of only 7 mrad. The detector noise is measured to be $\sigma_D = 0.0059 \pm 0.0005$. This corresponds at the points of most slope to a much larger interferometer phase noise of 69 ± 6 mrad. This was surprising as we expected the technical noise of the detector to generate roughly 6 mrad phase equivalent, as described in section 7.2.1. This seems to indicate some other source of noise, such as the frequency or intensity fluctuations in the imaging beams, that has yet to be confirmed.

The main source of noise at the points of large slope is the interferometer phase noise. This measurement found $\sigma_P(\Delta\Phi) = 0.0098 \pm 0.0006$, which corresponds to $\sigma_{\Delta\Phi} = 115 \pm 7$ mrad of interferometer phase noise. This agrees well with noise that was predicted in section 7.2.2 on the basis of the measured noise in the phase of the beat note. There, a single measurement with $2T = 50$ ms was predicted to have $\sigma_{\Delta\Phi} = 123 \pm 7$ mrad. This agreement indicates that the noise of the MEMS accelerometer is smaller than the acceleration noise of the atom interferometer, which is limited by the noise in the laser beat note phase. The benefit of using atom interferometer lies in its long term stability. After summing up all contributions at the steepest point of the fringe, we obtain an overall equivalent interferometer phase noise of $\sigma_{\Delta\Phi} = 134 \pm 7$ mrad, which, for this interferometer, amounts to an acceleration uncertainty of $16 \mu\text{m/s}^2$.

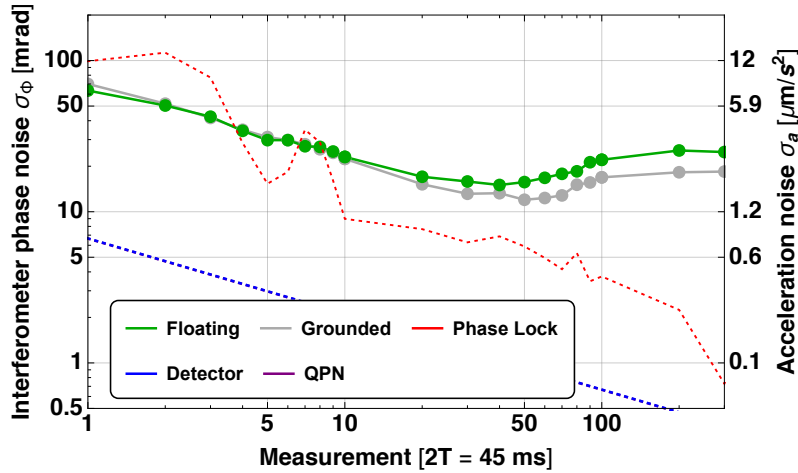


Figure 7.13: Allan deviations of the hybrid interferometer. Connected data points exhibit the residual noise after subtracting the accelerometer signal from the atomic signals. The red dashed line is the estimated Raman laser phase noise for a 45 ms interferometer. Blue and purple lines represent the noise levels due to the detector noise and the quantum projection noise, which overlap in the figure.

To assess how the acceleration noise of the atom interferometer behaves over long period of time, I extract the fit residuals of the population probability from Fig. 7.12 (b). These residuals are then converted into the interferometer phase noise and the accelera-

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tion noise at the steepest points of the fringe pattern. Figure 7.13 shows the Allan deviation of the converted interferometer phase as a function of the number of measurements in connected data points, where the acceleration noise is plotted on the auxiliary axis. The colour green and grey represent respectively the cases where the table is floated and not floated. As the acceleration has been taken out by fitting the interferometer fringes plotted against the MEMS accelerometer signal, the two residual interferometer phase noise overlap and agree with each other almost perfectly. This shows that the pendulum MEMS accelerometer has been hybridised with the atom interferometer to remove ambient vibrations down to the same noise level. In addition, the sensitivity of the interferometer appears to follow the phase lock noise of the Raman lasers, as shown in the red dashed line, which is converted into the interferometer phase noise analogously to Fig. 4.2 (b). As the number of measurements increases, we observe that the residual interferometer phase noise reaches the minimum after about 50 measurements, which currently correspond to 25 to 30 s of averaging. For averaging longer than 50 measurements, the interferometer phase noise deviates from the expected phase noise of the Raman laser beat note. The source of the additional noise is probably the intrinsic instability of the MEMS accelerometer due to temperature variations, as mentioned in section 7.2.3. In contrast the laser phase noise decreases continuously as the number of measurement increases, which shows the expected superior stability of the atom interferometer over the MEMS accelerometer.

7.3.4 Calibration

Accelerometer

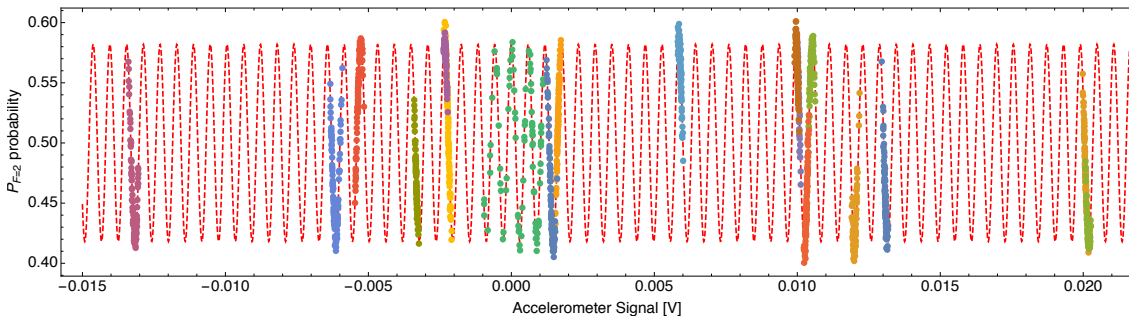


Figure 7.14: Acceleration measurements: 18 sets of data are combined together. The fit (dashed red) shows that the measured scaling factor of the pendulum accelerometer differs by 0.7 % from the datasheet.

Calibration of the atom interferometer depends on our knowing the laser wavelength and the time interval T , and the dominant uncertainty comes from T . I estimate that T is uncertain by 0.5 ns, making the scale factor uncertain by one part in 10^8 . This is

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insignificant compared with the uncertainty in the scale factor of the pendulum MEMS accelerometer. Hence, a long scan through the fringes can be used to calibrate the quartz pendulum accelerometer. We tilt the laser table by shifting weights in order to adjust the projection of local gravity g onto the Raman axis, and at each tilt angle, we make 200 measurements of the interferometer transition probability $P_{F=2}$ and the pendulum accelerometer voltage V . Figure 7.14 combines 18 sets of acceleration measurement at different tilt angles, indicated by different colours. The red dashed line is a single sinusoidal fit to all the data of the form

$$P = \frac{1}{2}(1 + C \cos(\alpha(V + V_0 + V_{\text{AccBias}}))), \quad (7.23)$$

with V_0 being an offset voltage and $\Delta\Phi_{\text{AccBias}} = \alpha V_{\text{AccB}} = 0.656321$ rad being determined using the bias voltage, specified by *Innalabs*. Using the relation $\alpha = k_{\text{eff}}T^2a/V$, we have determined the pendulum accelerometer calibration factor a/V to be

$$\frac{a}{V} = 1.3183 \pm 0.0001 \frac{\text{m/s}^2}{\text{V}}. \quad (7.24)$$

The manufacturer of the quartz pendulum accelerometer (*Innalabs*) specified 7939.85 $(\text{m/s}^2)/\text{A}$ at the room temperature. Using the measured resistance of our current monitor, 5981.1 ± 2.3 V/A , we therefore expected a calibration factor of 1.3275 ± 0.0005 $(\text{m/s}^2)/\text{V}$, which differs by 0.7% from Eq. 7.24. The specified temperature dependence of the calibration factor is only 30 ppm/ $^\circ\text{C}$, so that is too small to explain the discrepancy. As all specified values are tested in atmospheric pressure according to the data sheet, I wonder if the discrepancy is caused by our operating in ultra-high vacuum. We have not yet checked the offset voltage of the pendulum accelerometer. This could be done by comparing fringe patterns of different $2T$ -times: at zero acceleration, the phase of interference fringe pattern will be independent of T . Such measurement is scheduled in the further steps of the project.

Phase shifter

Until now, we have used the vibrations in the environment to scan the interference fringes. An alternative way is to vary the beat note phase on one of the three laser pulses, while fixing the phase of the other two phases. In the following experiment, I vary the beat note phase of the π -pulse, ϕ_2 . The black dots in Fig. 7.15 (a) show the $F = 2$ probability, as a function of the resulting interferometer phase $\Delta\Phi$. Here, the interferometer operates at $2T = 45$ ms on a floating table. We see that the interference fringe is quite blurred. In section 7.3.3, we expect the interferometer phase noise due to the noise in the beat

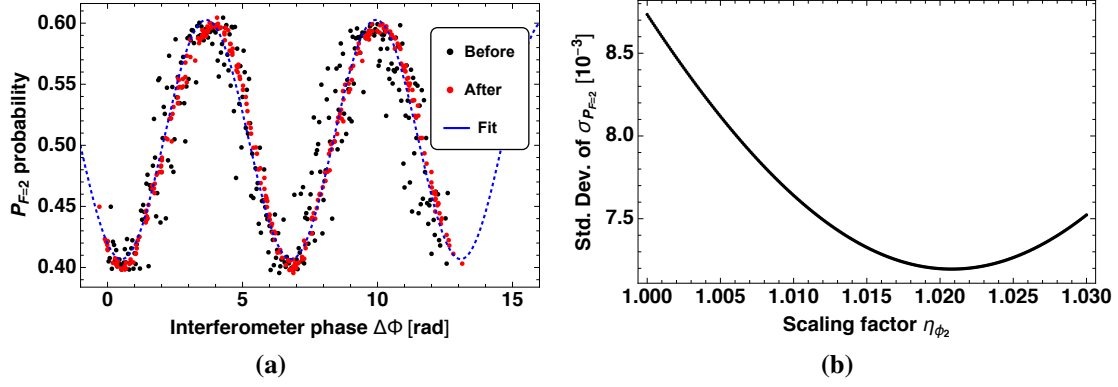


Figure 7.15: (a) Interferometer fringes scanned by varying the laser phase at the second Raman pulse ϕ_2 on a floating table. Black points: no vibration suppression. Red: vibration suppression using the Quartz pendulum accelerometer signal to re-bin the phase. (b) exhibits the optimisation curve of the scaling factor η_{ϕ_2} by minimising the standard deviation of the fit residuals.

note of the Raman laser to be around 100 mrad, which is significantly smaller than the noise in Fig. 7.15 (a). The blurred interference fringe in black is mainly due to the ambient acceleration noise on the floating table. To recover a clear interference fringe, the acceleration recorded by the MEMS accelerometer is first weighted by the triangular transfer function given in Eq. 7.11, then converted into an interferometer phase, $\Delta\Phi_{\text{MEMS}}$, using the calibration in Eq. 7.24 and finally subtracted from the applied phase $\Delta\Phi$. To allow the possibility of a mis-calibration of the phase shifter in the laser, which is set by a pre-calibrated analog channel, we have additionally introduced a scale factor η_{ϕ_2} . The interferometer phase is then taken to be

$$\Delta\Phi = -2\eta_{\phi_2} \phi_2 - \Delta\Phi_{\text{MEMS}}. \quad (7.25)$$

The red dots in Fig. 7.15 (a) show the same data after making these adjustments. Figure 7.15 (b) shows the standard deviation of the fringe fit residuals $\sigma_{P_{F=2}}$ as a function of η_{ϕ_2} . This has a minimum at $\eta_{\phi_2} = 1.0208$, which improves the interferometer phase noise by 20 %. With the optimised η_{ϕ_2} , we obtain a shot-to-shot interferometer phase noise of 70 mrad, which agrees well with the data in Fig. 7.13 after 1 measurement, where the fringes are scanned by vibrations.

7.4 Conclusion

In Chapter 7, we have started with the some theoretical treatment of the response of our atom interferometer to various noise sources, such as the noise in the phase of the beat note between the Raman lasers, ambient vibrations, quantum projection noise, shot

noise. Section 7.2 includes an experimental study of the contribution of the laser phase noise to the interferometer phase noise. We find that the *M-Squared* laser system contributes to the interferometer phase noise of around 124 mrad after every interferometer measurement, and that this noise goes down when the number of measurement increases. In section 7.3, we are able to take out the acceleration from the interferometer signal by plotting P_{F_2} against the MEMS accelerometer signal. By analysing the residuals of the fit of P_{F_2} , we understand that the main noise of the atom interferometer is limited by the phase noise of the Raman lasers in short times. As we average the interferometer phase over longer times, there is a strong indication that the noise is dominated by the temperature dependence of the pendulum MEMS interferometer. This emphasises the value of using atom interferometry to enhance the long term stability of the MEMS accelerometer. Furthermore, we explored the atom interferometer, by using its stability to calibrate the scaling factors of the MEMS accelerometer and the laser phase shifter.

Chapter 8

Summary and Outlook

8.1 Summary

The aim of the current on-going project in the Centre for Cold Matter at Imperial College London was to develop a compact, mobile 3-axis accelerometer using atom interferometry, which could be incorporated into an inertial navigation system. The first stage of the project is to prototype a working single-axis atom interferometer measuring the horizontal acceleration, perpendicular to the gravity, with a measurement noise close to $1 \mu\text{m/s}^2$. The success in the first stage is a result of the whole team working closely together in the field of engineering, modelling and programming. The focus of my PhD is on the hardware design and assembling for the project. With the help of my supervisor and post-docs, I have setup the whole experimental setup from an empty vacuum chamber and characterised and improved the performance of the device.

The highest, achieved data rate of our system has been demonstrated to exceed 4 Hz for an interferometer operating at $2T = 50 \text{ ms}$. With the available optical power of 60 mW in my 2D MOT, 10^8 atoms could be loaded into the 3D MOT in under 100 ms. The system has also shown the potential of even higher data rate with more optical power and further optimisations of parameters.

The 3D MOT is formed in the interferometer chamber by three pairs of MOT beams. With precise control of power, frequencies and polarisations, those beams were used in combination for launching atoms, polarisation gradient cooling, atomic state preparation and state detection. I showed that using a 4-step state preparation with optical and microwave pulses, around 94% of initially laser-cooled atoms could be concentrated in the state $|F = 1, m_F = 0\rangle$, ready for Raman light.

I made an improved version of Raman optics, which produced a collimated beam with a waist of 37 mm and was implemented in the UHV to avoid wavefront distortion, that would have been introduced by the viewports. Using this set of Raman optics, the

interference fringes that were first observed had very low contrast of approximately 10% for $2T = 5$ ms. We were able to measure the intensity of the Raman beam by studying the Rabi oscillation of the atoms and their optical pumping. The intensity was mapped by adjusting the launch trajectory of the atoms. These measurements indicated the presence of intensity lumps in the Raman beam. After avoiding these lumps, a factor of 3 of improvement of the fringe contrast was demonstrated.

I also started working towards a hybrid between the atom interferometer and a state-of-art Quartz pendulum accelerometer, where the atomic system is intended to correct for the bias drift of the MEMS device. The signal of the classical accelerometer was compared with signal of the atom interferometer. After optimising the integration between the two devices and subtracting the vibrations from the atomic system, we showed that the noise of the hybrid system is currently mainly limited by the noise in the phase lock in Raman laser system. This study enables us to distinguish different noise sources in the system and develop potential improvement for the next step of the project. Table 8.1 summarises the current key performance of our single-axis interferometer.

8.1.1 UK National Quantum Technologies Showcase 2018

As the first step towards operation in the field, the present single axis atom interferometer was presented at UK National Quantum Technologies Showcase in 2018 in collaboration with the laser company *MSquared*. Figure 8.1 (a) exhibits the sensing head without the enclosure on the left hand side of the picture; the laser and all control electronics are built into the device on the right hand side. Figure 8.1 (b) shows the system with the boxes closed. For the first time, our device was transported across the city of London and it was able to be set up within 3 hours to be operational. The interferometer operated continuously in the non-laboratory environment of the Queen Elizabeth II conference centre, with large fluctuations of magnetic field, temperature and acoustic noise. The sensing head was attached to a well-calibrated tilting stage. During the showcase, the stage was programmed to tilt the sensing axis by ± 1 mrad, which corresponds to a projected acceleration change between ± 1 mg. Figure 8.1 (c) exhibits the data collected at the showcase, where the signal of the atom interferometer (orange dots) closely follows the reference acceleration (black).

8.2 Outlook

Demonstrating the sensitivity to the horizontal acceleration in a noisy environment is only one of the first steps towards real-life operation of inertial sensors using atom inter-

	Expected	Measured
Vacuum		5×10^{-10} mbar
2D MOT		
Optical power		60 mW
Detuning (Cooling beam)		-12 MHz
Detuning (Push beam)		10 MHz
Loading rate		>10 Hz
3D MOT		
Optical power		90 mW
Detuning (Cooling beam)		-15 MHz
Quadrupole field gradient		15 G/cm
Atom number		10^8
Launch velocity	8.3 cm/s (H) 25 cm/s (V)	8.0 cm/s 23 cm/s
Raman beam		
Input beam power		1.5 W
Beam spot size	43 mm	38 mm
Beam collimation		≈ 0.3 mrad
Parallelism		≤ 17 μ rad
Mirror flatness		$\lambda/50$
Wave plate thickness uniformity		19 nm (PV)
Interferometer		
2T		45 ms
Contrast		0.15
Atom number		10^7 after 80 μ s selection
Interferometer phase noise (rms)		
Quantum projection noise	7 mrad	7 mrad
Detector noise	5 mrad	69 mrad
Laser phase noise	123 mrad	115 mrad
Acceleration noise		$1.5 \mu\text{m/s}^2$

Table 8.1: Summary of the key performance of the one-axis interferometer

ferometry. As pointed out in the thesis, further improvements are planned for the present system to increase the sensitivity of the interferometer. Firstly, the Raman optics will be replaced to avoid the intensity lumps. This will permit launching to extend the coherence time while still having good contrast. Secondly, the Raman laser system will be upgraded with a quieter phase lock system, which will lower the noise picked up by the interferometer. Thirdly, a new version of the light distribution unit is being developed using free-space optics, which has demonstrated higher polarisation purity and stability. This will improve the stability of the interference fringe contrast.

In addition to these, apart from scheduled system upgrades, we are planning a deeper

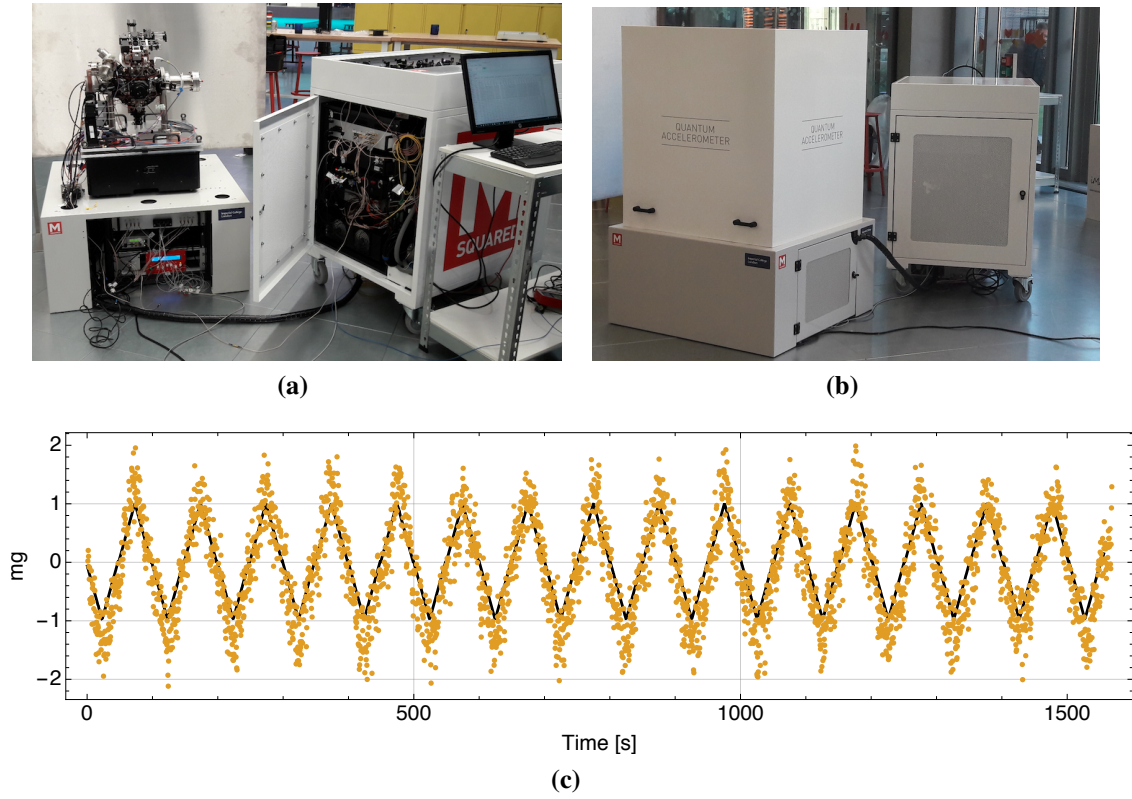


Figure 8.1: The atom interferometer with the laser system are shown without and with the enclosure in (a) and (b) respectively. (c) is a selection of data collected on the day of the showcase. The interferometer operated at $2T = 5$ ms and is able to track the reference acceleration (black) in a non-laboratory environment.

noise characterisation to distinguish the phase noise in the laser beat note from the MEMS accelerometer noise, and to improve the long-term stability of the interferometer together with the classical accelerometer. In particular, a study of the MEMS bias drift and its correction by the atom interferometer represents a crucial and necessary step towards enhanced INS, which will result in a significant reduction of the long-term positioning error. Another yet untouched aspect of developing a practical inertial sensor using atom interferometry is to distinguish and separate the interferometer phases due to acceleration and rotation. This work is the next research project to be conducted in the current 1-axis system.

In parallel, an extension of the present work has already been started, which involves building a universal single-chamber inertial sensor using atom interferometry. The new system will measure acceleration along both horizontal axes, and will then be extended to incorporate rotation measurements as well. Finally, I note that there is still more to be done from an engineering perspective to miniaturise and automate a commercial quantum sensor.

Chapter 9

Appendix

9.1 Double interferometer

As pointed out in section 2.1.3, Raman beams in retro-reflecting configuration enable two separate Raman transitions with opposite momentum transfer and therefore also sensitive to opposite velocities at a fixed Raman beat note frequency. This could lead to two simultaneously operating interferometers, which have been reported in Ref. [80]. Figure 9.1 shows the schematic of a double interferometers from left to right. The momentum dis-

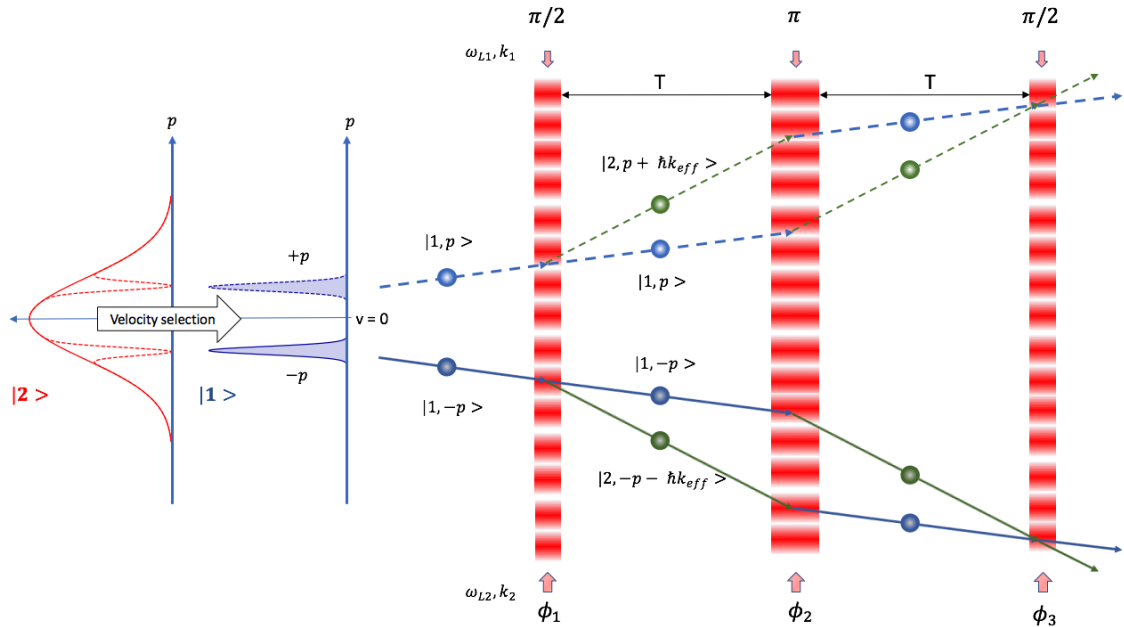


Figure 9.1: Double interferometers: Double interferometers occur with Raman beams in retro-reflection configurations. It arises when non-zero velocity selection is applied to a symmetrically distributed atom cloud centred at zero velocity. Figure motivated by Ref. [80].

tribution of atoms, which are prepared in $|2\rangle$, is illustrated by the red solid curve centred at zero velocity on the left of Fig. 9.1. If we choose the Raman beat note frequency to be

on resonance with atoms with a Doppler shift at the velocity $+v$ (or a momentum class of $+p$) for one pair of the Raman beams, the other one will be simultaneously resonant with atoms with a Doppler shift at the velocity $-v$. By applying a velocity selective pulse, we effectively select two separate atom clouds with opposite velocities, where the blue filled curves in Fig. 9.1 illustrate the selected momentum distributions. These two clouds undergo separate interferometer sequences simultaneously, which are illustrated by the solid and dashed interferometer paths.

Because of the finite cloud size and the small spatial separation between the selected atoms, it is not possible to resolve the two interfering atom clouds. Hence, the overall output after the interferometry is a sum of individual output of the interferometers, labeled by $P^+ = \frac{1}{2}(1 - \cos(\Delta\Phi^+))$ and $P^- = \frac{1}{2}(1 - \cos(\Delta\Phi^-))$, where superscript $\pm$ indicates interferometers with $\pm k_{\text{eff}}$ momentum transfers respectively. We write the total probability P^{tot} of atoms being detected on $F = 2$ as:

$$P^{\text{tot}} = P^+ + P^- = \frac{1}{2} \left(1 - \cos \left(\frac{\Delta\Phi^+ + \Delta\Phi^-}{2} \right) \cos \left(\frac{\Delta\Phi^+ - \Delta\Phi^-}{2} \right) \right). \quad (9.1)$$

$\frac{\Delta\Phi^+ + \Delta\Phi^-}{2}$ doesn't depend on the inertial forces. It, however, includes all non-inertial systematics. $\frac{\Delta\Phi^+ - \Delta\Phi^-}{2}$, on the other hand, contains only inertial effects, such as acceleration and rotation, whereas other systematics are canceled out. The cos term with the summed interferometer phases defines the fringe contrast of the double interferometers, which makes the acceleration measurement more vulnerable to systematic noise.

A straight-forward way to circumvent this issue is to suppress the contribution of the second cloud. Since it is impossible to single out only one pair of beams, we shift the velocity (momentum) distribution of the atom cloud in momentum space by giving it an initial velocity. In this case, the entire velocity distribution can be shifted away from zero velocity, causing an asymmetric distribution in the Raman spectrum. This, in turn, decreases the contribution of the interferometer one positive momentum transfer to the one with the negative momentum transfer, as the number of atoms resonant to the second interferometer decreases with increasing initial velocity. This is essentially also the motivation to launch atoms.

Figure 9.2 compares the fringe contrasts as a function of $2T$ time between interferometers operating on horizontally moving and resting molasses. The connected red data points show the contrast obtained from atom clouds with a horizontal launch velocity of $v_H = 0.1$ m/s along the Raman beams, whereas the black line shows the contrast without any horizontal launch velocity ($v_H = 0$). For the case with $v_H = 0$, a $12 \mu\text{s}$ velocity selection is performed at -70 kHz away from the atomic hyperfine transition frequency. Due to the retro-reflecting configuration of the Raman beams and the symmetric velocity

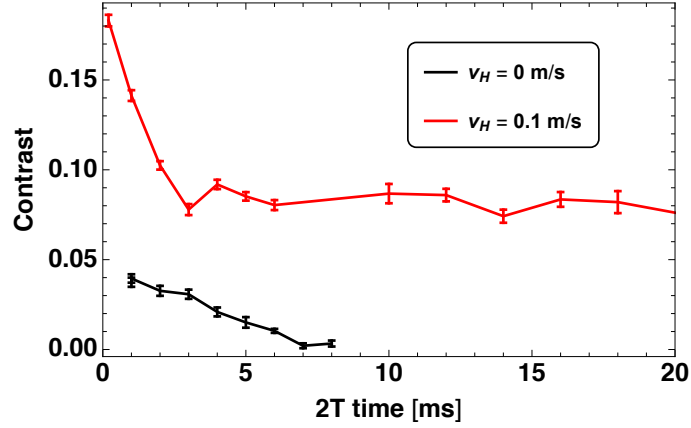


Figure 9.2: Interference fringe contrast plotted as a function of interferometer time $2T$. Velocity selection is applied on a resting atom cloud at ± 70 kHz away from the atomic resonance (black line), where the fringe contrast fades away completely after 7 ms of interferometer time. For moving clouds, a 40 velocity-selection pulse are performed on the Doppler resonance at -230 kHz (red line), exhibiting significantly stronger interference at much longer interferometer times.

distribution of atoms around zero, the second pair of Raman beams also select atoms, which are resonant with a detuning of $+70$ kHz, causing two competing interferometers working simultaneously. In addition to the Double interferometers, the interferometer pulses with a frequency width of around 75 kHz also off-resonantly drive the Doppler-insensitive transitions. All these contribute to a complete wash-out of the fringe pattern at $2T = 7$ ms. Hence, launching the atoms is necessary, as it prohibits the rise of Double interferometer and suppresses the off-resonant Doppler-insensitive transitions.

9.2 Noise in $|F = 1, m_F = \pm 1\rangle$

After the state preparation, there are approximately 6% of the initially laser-cooled atoms left in $|1, \pm 1\rangle$, which aren't affected by the interferometry, as described in section 3.8.2. They, however, contribute to the estimation of N_0 , which in turn affects the estimation of the fringe contrast. Let us write the total number of atoms after the velocity-selection pulse as $N_{\text{tot}} = N_0 + N_{\pm}$, where the number of atoms in $|1, \pm 1\rangle$ is denoted by $N_{\pm} = \epsilon N_{\text{tot}}$ and ϵ is the proportional factor. We express the total number of laser-cooled atoms N_{tot} as:

$$N_{\text{tot}} = N_0 + N_{\pm} = N_0 + \frac{\epsilon}{1 - \epsilon} N_0 = \frac{1}{1 - \epsilon} N_0. \quad (9.2)$$

Using this, we write the $F = 2$ probability before correcting for the residual atoms, $\tilde{P}_{F=2}$, as:

$$\tilde{P}_{F=2} = \frac{N_2}{N_{\text{tot}}} = (1 - \epsilon) P_{F=2}, \quad (9.3)$$

where $P_{F=2} = N_2/N_0$ is the corrected $F = 2$ probability. Analogously, the uncorrected contrast $\tilde{C}$ is given by:

$$\tilde{C} = (1 - \epsilon) C. \quad (9.4)$$

Hence, the interferometer phase noise $\tilde{\sigma}_{\Delta\Phi}$ before correcting $N_{\pm}$ contribution is expressed in terms of $P_{F=2}$ and N_0 using Eq. 9.3 by:

$$\begin{aligned} \tilde{\sigma}_{\Delta\Phi} &= \frac{2}{\tilde{C}} \sigma_{\tilde{P}_{F=2}} = \frac{2}{\tilde{C}} \sqrt{\left(\frac{d\tilde{P}_{F=2}}{dP_{F=2}} \right)^2 \sigma_{P_{F=2}}^2 + \left(\frac{d\tilde{P}_{F=2}}{d\epsilon} \right)^2 \sigma_{\epsilon}^2} \\ &= \frac{2}{C} \sqrt{\sigma_{\text{QPN}}^2 + P_{F=2}^2 \frac{\sigma_{\epsilon}^2}{(1 - \epsilon)^2}} \\ &\approx \frac{2}{C} \sigma_{\text{QPN}} \left(1 + \frac{P_{F=2}^2}{(1 - \epsilon)^2} \frac{\sigma_{\epsilon}^2}{\sigma_{\text{QPN}}^2} \right), \end{aligned} \quad (9.5)$$

where σ_{ϵ} is the standard deviation of ϵ and $\sigma_{P_{F=2}}$ is assumed to be σ_{QPN} described in section 7.1.3. The approximation is due to the Taylor expansion about $\sigma_{\epsilon} = 0$. Equation 9.5 shows that to lower $\tilde{\sigma}_{\Delta\Phi}$, it is more important to stabilise $N_{\pm}$ than to minimise it.

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