Laser Cooling and Trapping of Metastable Helium Atoms

A thesis submitted for the degree of Doctor of Philosophy of the Australian National University

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Statement of Authorship

This thesis is an account of research undertaken in the Atomic and Molecular Physics Laboratories and Laser Physics Centre, Research School of Physical Sciences and Engineering, Australian National University between February 1995 and August 1998.

Except where acknowledged in the customary manner, the material presented in this thesis is, to the best of the author’s knowledge and belief, original and has not been submitted for a degree at any university.

Dragana V. Jablan Milic

February, 1999
Mami i tati.
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The staff and students of AMPL in late 1997. Back row: Dr Anatoli Kheifets, Dr Malcolm Elford, Professor John Carver, Professor Erich Weigold, Dr Brenton Lewis.

Third row: Dr Stephen Buckman, Stephen Battisson, Kevin Roberts, Jennie Gibson, Dr Stephen Gibson, Weijian Liu, Dr Alan Lun, Professor Robert Crompton.

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Front row: Graeme Cornish, James Sullivan, Dr Maarten Hoogerland, Kevin Lonsdale, Dr Julian Lower.

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Summary

Intense, highly collimated sources of atoms have many potential applications, such as the study of atom optical devices, e.g., for loading atoms into hollow fibre waveguides. In atomic physics, many collision processes can only be measured with the sensitivity offered by such high flux sources. With the advent of laser cooling mechanisms, new ways have been developed to create ultracold, ‘ultra-bright’ atomic beams in a background free environment. Some of this concepts are introduced in chapter 1. In chapter 2 the basics of some of the laser cooling and trapping mechanisms are outlined. In chapter 3 the fundamentals of some ‘sub-Doppler’ cooling mechanisms are explored experimentally, using sodium atoms.

An important part of this thesis is the successful construction of a bright source of metastable (2^3S) helium atoms. A standard, liquid nitrogen cooled, DC discharge source produces an atomic beam with a flux of \(10^{14} - 10^{15} \text{ s}^{-1} \text{sr}^{-1}\). The beam is transversely collimated using effectively curved wavefronts, yielding a capture angle of 0.005 sr. It is then slowed from its initial axial velocity to 100 ms\(^{-1}\) employing a Zeeman tuned, counterpropagating laser beam. It is then focussed on a 3 mm aperture and subsequently recollimated. The bright atomic beam is currently ready to be used for electron-atom scattering experiments. In Chapter 4 the operation of each of the beam brightening stages is described in detail. All of these cooling stages employ diode lasers to excite the \(2^3S_1 \rightarrow 2^3P_2\) transition. Details on the diode laser drivers, and further electronics used for frequency control of the lasers, can be found in Chapter A.

Using different settings for the various laser and magnetic fields, the bright beam can be used to load a magneto-optic trap, where the atoms are subsequently cooled to temperatures in the mK range. A trapping chamber has been constructed, and preliminary experiments are presented in Chapter 6 of this thesis.

These experiments use a newly developed laser source of UV radiation, resonant with a higher transition in helium. The design of the laser source, a frequency
doubled titanium:sapphire laser, is detailed in chapter 5.
Contents

1 Introduction
   1.1 Atom Optics ................................................................. 1
   1.2 Metastable Helium ....................................................... 1
       1.2.1 Atomic and electronic collision experiments .......... 2
       1.2.2 Atom optics experiments ..................................... 2
       1.2.3 The Bright Source ............................................. 4
   1.3 The scope of this thesis ............................................. 5
   1.4 Contribution of others ............................................. 6

2 Light forces
   2.1 Radiation pressure ................................................... 8
   2.2 Doppler Cooling ....................................................... 9
   2.3 Sub-Doppler Cooling ............................................... 12
   2.4 VSCPT .................................................................. 18

3 Transverse laser cooling of a velocity-selected sodium atomic beam
   3.1 Introduction ............................................................ 27
   3.2 Experiment ............................................................ 32
   3.3 Results ................................................................. 38
   3.4 Discussion ............................................................ 41
   3.5 Conclusion ............................................................ 43

4 The bright beam apparatus
   4.1 Introduction ............................................................ 50
   4.2 LN$_2$ Cooled metastable atom source ....................... 51
   4.3 Collimation ............................................................ 54
   4.4 Slowing ................................................................. 57

viii
CONTENTS

4.4.1 Design ................................................................. 57
4.4.2 Deceleration results .............................................. 61
4.5 Compression ............................................................. 63
4.5.1 Introduction .......................................................... 63
4.5.2 Simulation ............................................................ 66
4.5.3 Design ................................................................. 67
4.5.4 The quadrupole magnetic field .............................. 68
4.5.5 Operation ............................................................ 70
4.5.6 Conclusions and suggested improvements ............ 71

5 Generation of 389 nm light ........................................ 74
5.1 Motivation ................................................................. 74
5.2 Frequency doubling .................................................. 75
5.3 Setup ....................................................................... 78
5.4 The Crystal ............................................................... 80
5.5 The Cavity ................................................................. 82
  5.5.1 The Cavity Type .................................................. 83
  5.5.2 The Cavity Design .............................................. 84
5.6 The Locking Scheme ................................................ 86
5.7 Implementation ....................................................... 88
5.8 Conclusions ............................................................ 92

6 Trap experiments ....................................................... 94
6.1 Introduction ............................................................. 94
6.2 The design of the trap .............................................. 98
6.3 Obtaining the MOT .................................................. 99
6.4 Cold Collisions Experiments ................................. 101
6.5 Cold Collision Results and Discussion ................. 103
6.6 Conclusions and suggested improvements ............ 106

7 Conclusion ............................................................... 108
7.1 Future research ...................................................... 109

A Electronic devices developed for the bright beam machine and related experiments .................................. 111
A.1 Diode laser stabilisation ........................................... 111
A.2 Diode laser controllers ........................................... 112
CONTENTS

A.2.1 Circuit description ...................................................... 112
A.2.2 Noise reduction and stabilised power supply ..................... 119
A.2.3 Lock-in amplifier ............................................................ 120
A.2.4 Results ................................................................. 122
A.3 AOM drivers ................................................................. 123
   A.3.1 Digital to analog converter (DAC) summing amplifiers .... 124
   A.3.2 VCO frequency comparator ........................................ 125
   A.3.3 VCO-RF board ......................................................... 125
Chapter 1

Introduction

1.1 Atom Optics

The concept of cooling atoms with laser light to low temperatures was first proposed by Hänsch and Schawlow [1] and Wineland and Dehmelt [2] in 1975. Since the first demonstration of one dimensional laser cooling in 1981 [3], two dimensional cooling in 1982 [4] and finally three dimensional laser cooling [5] this field has grown rapidly, and last year's (1997) Nobel Prize in Physics was awarded to Prof. C. Cohen-Tannoudji, Dr. W.D. Phillips and Prof. S. Chu, for their groundbreaking work in this field. The achievement of Bose-Einstein condensation in 1995 [6] was one of the most exciting breakthroughs in modern atomic physics, with laser cooling as an essential part of the process for forming the condensate. Laser cooling has led to the field of atom optics, a young discipline [7] which involves studying ways to manipulate neutral atoms in a way similar to the way "classical" optical elements are used to manipulate light. Ultimately, atom optics studies the properties of neutral atoms which are so cold that their de Broglie wavelength $\lambda_{dB} = \frac{h}{p}$ is comparable to the wavelength of the laser fields used to manipulate them.

1.2 Metastable Helium

The work described in this thesis focuses on the development of an apparatus for production of a "bright" beam of metastable (long lived) ($2^3S$) helium atoms using laser cooling techniques. We have used the 1083 nm transition ($2^3S_1 \rightarrow 2^3P_{0,1,2}$) as the cooling transition. The lower, or ground state for this transition
is the metastable $^{23}\text{S}_1$ state which is the lowest excited level in atomic He. It is metastable because decay to the ground state is forbidden by the dipole selection rules on both spin (triplet $\rightarrow$ singlet) and angular momentum $l = 0 \rightarrow l = 0$. However it does have a finite decay width to the true He ground state ($^{1}\text{S}_0$) with an average decay time of approximately 2 hours, as calculated by Drake [8], and partly confirmed by the measurements by Woodworth and Moss [9]. The upper $^{23}\text{P}_{0,1,2}$ states can also decay, via a magnetic quadrupole interaction, to the $^{1}\text{S}_0$ state, with a mean decay time of approximately 0.1–1 s, which is not a significant problem for the present work, as both the flight time and the lifetime in the trap are significantly shorter than this. The energy levels and lifetimes of states for the lowest lying states in helium are shown in figure 1.1.

In the present work metastable helium atoms are interesting for two classes of experiments: Atomic and electronic collision experiments and atom optics experiments.

1.2.1 Atomic and electronic collision experiments

The simplicity of the helium atom makes it a good model system for calculations of atomic and electronic collision processes, and hence experiments in these areas on metastable helium are required to discriminate between the theoretical models used.

Metastable atoms are important in atmospheric, gas discharge and laser physics [10, 11]. Their long lifetime means that their influence on the dynamics of, for example a gas discharge, can be significantly larger than their density alone may imply.

1.2.2 Atom optics experiments

For light He atoms the velocity change $v_{\text{rec}}$ arising from the recoil due to the absorption of a single photon for the 1083 nm transition is very large (see table 2.1). This is even more so for the 389 nm transition to the $3^3P$ states (figure 1.1). The large recoil should allow the creation of large angle atomic beamsplitters, and a detailed study of the effects of the quantisation of the atomic motion in a laser field.

Metastables de-excite on collision with surfaces, by transferring their energy to a surface electron. As this energy (20 eV for metastable helium) is much
larger than the work function of the material, an electron is ejected from the surface. This enables the detection of only those atoms which interact with laser light forces. This is of importance for tests of many atom-optical devices, such as atomic mirrors and waveguides, where light forces repel the atoms from the surface. Another consequence is that the emitted electrons can be detected with near unity detection efficiency by channeltron or multi-channelplate detectors.

Their 20 eV excitation energy may be used to efficiently damage a layer of “photoresist”, formed on a thin layer of gold on a silicon wafer. In preliminary experiments, a photoresist formed by a “Self Assembled Monolayer” (SAM) of dodecanethiolate was exposed to the metastable helium beam through a transmission mask. The mask pattern was reproduced with sub-micron accuracy, and with an excellent contrast in the silicon [12]. Traditional techniques use UV light for the exposure, and the resolution is limited by the wavelength of the light.
Due to the short deBroglie wavelength of a massive particle, such as a helium atom, and the possibility to manipulate the atoms with laser light, this technique holds some promise for semiconductor nanofabrication.

1.2.3 The Bright Source

The above experiments have requirements with regard to flux, brightness and collimation of the He* beam, which are much higher than is achievable with a "standard" nozzle source for He*. This is the prime motivation for the development of the bright beam machine described below.

This thesis reports the successful construction of a bright source of metastable (2S) helium atoms (chapter 4). The liquid nitrogen cooled, DC discharge source produces an atomic beam with a flux of $10^{14}$-$10^{15}$ s$^{-1}$sr$^{-1}$. The beam is transversely collimated using effectively curved wavefronts, yielding a capture angle of 0.005 sr. It is subsequently slowed from its initial axial velocity to $\sim$100 ms$^{-1}$ employing a Zeeman tuned, counterpropagating laser beam. It is then focused to a $\sim$3 mm diameter spot and recollimated once again. We use diode lasers (SDL 6701-H1) to excite the $2^3S_1 \rightarrow 2^3P_{0,1,2}$ transitions for all beam brightening stages and experiments.

The techniques used to create this "bright beam" are similar to those used to create bright beams of sodium [13], neon [14, 15] and helium [16] modified to increase the capture range and optimise the yield of the brightening process for metastable helium.

Other methods to create a bright, slow atomic beam, such as the 'LVIS' scheme [17] are not suitable for metastable atoms, as they depend on a finite vapour pressure of the atoms in a vapour cell. Metastable atoms however de-excite on a collision with the vapour cell wall, and hence it is not feasible to have a finite 'vapour' pressure of metastable atoms.

In several other laboratories metastable atoms are trapped in a magneto-optic trap, in a way similar to the way described in chapter 6 [18, 19, 20]. However, for many experiments, particular in atomic physics, it is desirable to have a continuous, monochromatic beam of slow atoms, as is described here.
1.3. **THE SCOPE OF THIS THESIS**

### 1.3 The scope of this thesis

Chapter 2 of this thesis is an overview of the concepts and theory of laser cooling. These concepts are used throughout the remainder of this thesis. For a more in-depth review the reader is referred to [21], [22], [23], [24]. As a final section of this chapter a short review of "velocity selective Population Trapping" is included, as this is an important subject for future research. Chapter 3 describes a detailed study of one dimensional, polarisation gradient laser cooling of sodium atoms. The limits to polarisation gradient cooling of the transverse atomic velocity are of considerable interest, both theoretically and experimentally. Central to this interest is the dependence of the transverse atomic temperature on the laser intensity. An extensive study of the transverse velocity distribution of polarisation gradient cooled sodium atoms for the \( \{3s\}^2S_{1/2} \ (F=2) \rightarrow \{3p\}^2P_{3/2} \ (F=3) \) transition at 589.1 nm was conducted. A detailed analysis of the experimental results provides a better insight into the efficacy of the polarisation gradient cooling process in regard to cooling laser intensity, detuning and polarisation. Sub-recoil resolution measurements have been performed on the transverse cooling of a velocity selected sodium atomic beam in a transverse laser field, consisting of two, counterpropagating, laser fields with orthogonal linear polarisation. The insights acquired in these experiments have been used in the design of the post-compression collimation stage of our bright beam source, to further narrow the transverse spread of the already compressed slow beam of helium atoms.

A detailed description of the bright beam apparatus can be found in chapter 4 of this thesis. The emphasis lies on the development of the final stages of the beam brightening apparatus: the compression and post-collimation sections. The description of the metastable atom source, and the collimating and slower sections are included for completeness only, as these have been mainly set up by my coworkers (see section 1.4).

An important aspect of this thesis was the design of the driving, locking and stabilisation electronics for the diode laser systems, as well as the drivers for the Acousto-Optic Modulators (AOMs) used in the trap experiments. Details on these designs can be found in chapter A. The diode laser electronics are used in every stage of the 'bright beam' and for the metastable atom trap. The design of a number of detectors, proportional-integral-differential (PID) controllers and timing electronics circuits are more standard and are hence not included in this thesis.
CHAPTER 1. INTRODUCTION

In parallel with the development of the bright beam machine I have developed a laser source for 389 nm UV light, corresponding to the $2^3S_1 \rightarrow 3^3P_{(0,1,2)}$ transitions in helium, which has a recoil velocity $v_{\text{rec}} = 25.6 \text{ cms}^{-1}$, which is three times bigger than that for the 1083 transition. To do this, the output of a 778 nm Ti:Sapphire laser beam is frequency doubled in an LBO crystal in an external build up cavity. Chapter 5 describes the steps in development of the 389 nm blue light in detail.

Finally, we have developed a Magneto-Optical-Trap (MOT), which is loaded from the collimating and cooling stages of the bright beam source. While trapping metastable He atoms using the infrared 1083 nm transition, we have studied the change of the ionising collision rate whilst probing with a different colour laser (389 nm), in the absence of the trapping laser fields. The description of the trap and the collision dynamics experiment in the trap are given in chapter 6 of this thesis.

1.4 Contribution of others

The development of such a complicated and sophisticated apparatus as our bright beam source required the input of more than one person and I would like to mention the contributions of my fellow students, supervisors, and technical staff.

- The first stage of our apparatus - A liquid nitrogen cooled source of metastable He atoms was designed and put together by Mr Stephen Battison, of the technical staff in the Atomic and Molecular Physics Laboratories. I have done initial calculations prior to the design of the source and took part in the designing process and setting up of the source. The measurements for the source characterisation were performed by my student colleague Weijian Lu.

- The collimation stage of our apparatus was designed following a similar device used for metastable neon in the Technical University of Eindhoven, the Netherlands, and was put together, aligned and tested by Mr Roel Knops, a visiting student from the same University. My input to the development of this stage of the apparatus was the calculation of the incident angle of the laser beam entering the collimating stage, which would result in the correct dependence of the angle between the laser beam and the atomic beam as a
1.4. CONTRIBUTION OF OTHERS

function of position. The design work was done by Mr. Ian McRae of the technical staff in the Laser Physics Centre.

• The development of the Zeeman slower was primarily the task of my student colleague Weijian Lu [25]. The design of the hardware and the winding of the coils was done by Mr Graeme Cornish of the AMPL technical staff.

• The development of the compression stage was primarily my task and Mr Lindsay Wilson, of the AMPL technical staff has made a significant contribution to the design of the optical mounts in this stage. Following this, Mr Wilson has, as the permanent technical staff member on our team, made many contributions to the experiments described in this thesis.

• My student colleague Massimiliano Colla has put together the hardware of the MOT and carried out the initial alignment of the trap, prior to its attachment to the rest of the apparatus. He has, as well made a contribution to my final "trap experiment", see chapter 6 of this thesis. The trap chamber was designed by Mr Graeme Cornish.
Chapter 2

Light forces

2.1 Radiation pressure

Resonant laser light can be used to exert a force on neutral atoms. An atom that interacts with a resonant laser field of wavelength \( \lambda \), can absorb a photon of laser light (figure 2.1). At the same time, the atom will absorb the photon momentum \( \hbar \vec{k} \). Here \( \vec{k} \) is the wave vector of the photon:

\[
|\vec{k}| = \frac{2\pi}{\lambda}
\]  

(2.1)

Figure 2.1: (a) Atom is excited by absorption of a laser photon (b) At the same time, the atom absorbs the photon momentum. (c) After spontaneous emission of a photon in a random direction, the atom returns to the ground state, and can absorb another laser photon. After many of these cycles, the recoils due to spontaneous emission average to zero, and the atom is accelerated in the direction of the laser.
2.2. DOPPLER COOLING

The excited atom has changed its velocity due to this single recoil by an amount \( v_{\text{rec}} = \frac{\hbar k}{m} \). This is the definition of the recoil velocity, and it is the consequence of momentum conservation. After the natural lifetime the atom spontaneously emits a photon in a random direction (figure 2.1c), receiving a recoil in the direction opposite to that of the spontaneously emitted photon. This brings the atom back to the ground state, from which it can be excited again. After a number of these cycles the resultant recoil from the spontaneous emissions will average to zero, while absorption of the laser photon’s recoil gives the atom a momentum change in the direction of the laser light. If the laser light is resonant with a strong electronic transition in the atom, the atom can absorb many photons per unit time. The maximum spontaneous emission rate is \( \Gamma/2 \), with \( \Gamma = 1/\tau \) the reciprocal of the natural radiative lifetime of an atom in the excited state. An atom can spend, at most, half of its life in the excited state. The maximum force exerted on the atom is:

\[
F = \frac{\hbar k \Gamma}{2},
\]

(2.2)

For helium, excited on the \( 2^3S_1 \rightarrow 2^3P_2 \) transition at 1083 nm, this force is \( F = 3 \times 10^{-21} \) N, corresponding to an acceleration of \( a = 4.6 \times 10^5 \) ms\(^{-2} \). This large acceleration is the key to the use of laser radiation pressure to control the motion of atoms in our bright beam machine.

### 2.2 Doppler Cooling

We assume that the atom is a two level system having only a ground \( |g\) and excited \( |e\) state (figure 2.3(a)), with a two-level atomic transition frequency \( \omega_a \). The laser has a frequency \( \omega_L < \omega_a \). For a moving atom, a laser that is not exactly on resonance and which has a finite laser power, the force in equation 2.2 is modified by the Lorentzian line shape of the transition, the intensity of the laser light and the Doppler shift \( \Delta \omega_D = -\vec{k} \cdot \vec{v} \):

\[
\vec{F}_i = \frac{\Gamma}{2} \frac{s}{1 + s + 4 \left[ \frac{\Delta \omega_D}{\Gamma} \right]^2} \hbar \vec{k}_i
\]

(2.3)

where \( \vec{v} \) is the velocity of the atom, \( s = I/I_0 \) is the saturation parameter, defined as the ratio between the laser intensity \( I \) and the saturation intensity \( I_0 \). \( \Delta = \omega_L - \omega_a \) is the laser detuning. The saturation intensity is given by [26]

\[
I_0 = \frac{\pi \hbar c \Gamma}{3 \lambda^3}
\]

(2.4)
In a counterpropagating laser fields configuration (figure 2.3) using slightly red detuned ($\Delta < 0$) laser light, each of the laser beams forming the standing wave exerts radiation pressure on the atom. Since the average force from the spontaneously emitted photons vanishes, as they are emitted randomly, the net force from each of the beams is in the direction of that beam and is given as the product of the photon scattering rate and the momentum change $\hbar k_i$ per scattering event.

Figure 2.3: For cooling in one dimension, counterpropagating laser fields are used.

When an atom moves it ‘sees’ the laser beam Doppler shifted by an amount $\Delta \omega_D = -kv$. That effectively means that the counterpropagating laser beam will appear blue shifted to the atom, and therefore closer to the resonance. As we can see from the equation 2.3 the term in brackets is velocity dependent, and is a function of the Doppler shift $-k_i \cdot \vec{v}$ and laser detuning $\Delta$, with a maximum at $v = \Delta/k$ (for which the Doppler shift brings the laser into resonance) and a FWHM (full width half maximum) equal to $\Gamma/k$. Assuming a small laser intensity $I \ll I_0$ we can add the two laser forces independently, producing the situation shown in figure 2.4. Near zero velocity, the atom experiences a slowing force proportional to its velocity, as if it were moving in molasses (hence the term ‘optical molasses’. Thus the motion of the atom is strongly damped with a friction coefficient $\alpha$ that is linear in the laser intensity. The width of the velocity distribution is thus greatly narrowed. If the velocity distribution is Gaussian, an equivalent temperature:

$$\frac{1}{2}T = \frac{1}{2}\frac{m \langle v^2 \rangle}{k_B}$$

(2.5)

can be assigned to the atomic ensemble. As the velocity distribution is narrowed, this temperature decreases, and is hence known as laser cooling. By this mechanism alone, the atoms could be cooled arbitrarily close to absolute zero.
2.2. DOPPLER COOLING

Figure 2.4: a) Radiation pressure force b) Fluctuations of the radiation pressure force due to random recoils from the spontaneously emitted photons that heat the atom ensemble, imposing in this way a limit to laser cooling. The fluctuations are simulated here by adding a random number, proportional to the heating rate, to the force (see equation (2.8)).

However, the force has fluctuations around the curve shown in figure 2.4(a), due to the random nature of spontaneously emitted photons, which can be illustrated as in figure 2.4(b). This leads to a finite ultimate temperature, as detailed below.

The random recoils from the spontaneously emitted photons heat the atom ensemble, imposing in this way a limit to laser cooling. Fluctuations of the optical force around its average value originate from the discrete transfer of momentum, which generates this force in the first place. Force fluctuations are due both to the randomness in the direction of the spontaneously emitted photons, and to the the randomness in the absorption of an emitted photon. The atomic motion in a momentum space can be thought of as random walk process, in steps of recoil $\hbar k$ of an atom. The mean square momentum after N cycles is given by:

$$\langle p_z \rangle^2 = 2N\hbar^2k^2$$  \hspace{1cm} (2.6)

This can be described with a diffusion coefficient in velocity space $D$. The two processes are in balance when the velocity distribution of the atoms reaches steady state. By integrating the Focker-Planck equation [7] for this ensemble, an equilibrium Gaussian distribution is obtained, meaning that the atoms may be characterised by a temperature. The cooling rate can be found by using the friction
coefficient $\alpha$:

$$\left( \frac{dE}{dt} \right)_{\text{cool}} = \frac{d(\frac{1}{2}mv^2)}{dt} = \frac{vd(mv)}{dt} = F_v = -\alpha v^2 = \frac{2\alpha E}{m} \quad (2.7)$$

The heating rate is calculated using the assumption that at each spontaneous emission event, energy corresponding to the recoil of one photon with momentum $\hbar k$ is added to the system. This gives a heating rate:

$$\left( \frac{dE}{dt} \right)_{\text{heat}} = 2 \cdot \frac{1}{2} m \left( \frac{\hbar k}{m} \right)^2 \frac{\Gamma s}{2(1 + 4\Delta^2/\Gamma^2)} = \frac{\hbar^2 k^2 \Gamma s}{4m} \equiv \frac{2D}{m} \quad (2.8)$$

In the last step we have assumed a detuning $\Delta = -\Gamma/2$, which achieves optimal cooling in this scheme. The first factor of two is due to the fact that we have two laser beams. Fluctuations in the force give rise to another factor of two in the heating rate. In equilibrium we end up with:

$$\frac{2\hbar k^2 s}{m} E = \frac{\hbar^2 k^2 \Gamma s}{2m} \Rightarrow E = \frac{D}{\alpha} = \frac{\hbar \Gamma}{4} \quad (2.9)$$

This result is known as the Doppler limit. This result is surprising, as it seems that the equilibrium kinetic energy does not depend on the mass of the atom, the wavelength of the transition or the laser intensity. It only depends on the linewidth of the transition used for the cooling. As (in equilibrium) the velocity distribution is Gaussian, an 'equivalent temperature' $\frac{1}{2}kT = E$ (in one dimension) can be assigned to an ensemble of atoms. For sodium, $T = 240 \ \mu K$. The one dimensional Doppler limit for metastable helium is $T_D = 39 \mu K$. This cooling scheme can easily be generalised to two or three dimensions, as illustrated in figure 2.5. Depending on whether the laser cooling is performed in one, two or three dimensions (figures 2.3 and 2.5(a) and (b)), the Doppler limit has a different value, as each standing wave damps one velocity component.

### 2.3 Sub-Doppler Cooling

Although some of the first measurements confirmed the Doppler limit [5], in many experiments much lower temperatures were observed [27]. These lower temperatures were explained by subsequent theoretical developments [24]. Here I give an overview of the basic concepts.

The electronic levels of an atom in a strong (resonant) light field are shifted by the AC Stark shift equal to $\hbar \Omega_s = \sqrt{\Omega^2 + \Delta^2} - \Delta$, with $\Omega = \vec{d} \cdot \vec{E}/\hbar$ the
2.3. **SUB-DOPPLER COOLING**

(a) Two dimensional cooling  
(b) Three dimensional cooling: "optical molasses"

**Figure 2.5: Two- and three-dimensional Doppler cooling setups.**

On resonance Rabi frequency, where $\vec{d}$ is the electric dipole and $\vec{E}$ is the electric field vector. For many practical situations ($\Delta \gg \Gamma, \Omega$) the light shift can be approximated as:

$$\hbar \Omega_s = \frac{\hbar \delta \Gamma^2}{8 \Delta} = \frac{\hbar \Omega^2}{4 \Delta}$$

In this limit, the light shift is proportional to the laser intensity and inversely proportional to the laser detuning. Hence, a gradient in laser intensity and thus in the AC Stark shift will exert a force $F = -\nabla \Omega_s$ on the atoms. Large gradients in laser intensity occur, for instance, in evanescent waves formed by total internal reflection of a laser beam in a prism. In a standing wave, there are large gradients in the light shift between the nodes and the antinodes and thus the atoms will experience a force, apart from the Doppler force. If the atom has more than two levels, e.g., in the case of magnetic degeneracy of the excited and ground states, these forces can be used to cool the atoms. The coupling of the magnetic sub-levels to the excited state by polarised laser light is determined by the Clebsch-Gordan coefficients for the transition matrix elements. Hence a spatial gradient of the polarization of the light can also give rise to gradient forces.

We will first consider an atom in a laser field generated by the superposition
of two counterpropagating laser fields, linearly polarized at right angles to one another (lin-perp-lin). The polarization of the resulting laser field has a gradient: the polarization changes in space such that every \( \lambda/8 \) (see Appendix Chapter 2 figure 2.11), we have a different polarization. This is the key to understanding the Sub-Doppler cooling process.

The above polarization fields labeling, which is the same as the one used in figures 2.9 - 2.12 denotes the polarization of the two counterpropagating laser fields along \(-z\) and \(+z\) direction respectively, where "\( \pi \)" or "lin" stands for linearly polarized light and "\( \sigma \)" for circularly polarised light. The sense of polarization is denoted by the superscripts, with \( \pi^\pm \) for the \( E \) polarised along the \( x \)-axis, and \( \sigma^\pm \) for the \( E \) angular momentum in the \(-z\) direction. In geometrical optics, circularly polarised light is labeled according to its helicity. In atom optics circularly polarised light is labeled according to the projection of its angular momentum along the appropriate (in our case \( z \)) axis. The selection rule for emission and absorption of \( \sigma^\pm \) photons: \( \Delta m = \pm 1 \), is in accordance with labeling the polarization as a function of angular momentum. The photon angular momentum is conserved after the reflection of a photon at an incidence angle of 90°, which is not the case with helicity. In a setup like the one to be described here, with a large number of laser beams in different directions, this way of polarization field labeling is very convenient.

The different ground state sublevels in an atom have different light shifts due to their different Clebsch-Gordan coefficients. Figure 2.6(a) shows a level scheme for an atom with a \( J_g = \frac{1}{2} \rightarrow J_e = \frac{3}{2} \) transition and the resulting light shift energies, \( U = \hbar \omega \), for the two ground states. In \( \sigma^+ \) polarized laser light, that induces \( \Delta m = 1 \) transitions, the \( m = +\frac{1}{2} \) magnetic substate is light shifted the most: its coupling with the excited state is the largest. When the laser detuning is negative, this means that this substate is the lowest in energy. In other words the light shift is negative and the optical pumping at any given position tends to drive the atom to the state with the lower potential. After many absorption-spontaneous emission cycles, the atoms are optically pumped to this state. We will now consider an atom that is in this particular ground state, in \( |m\rangle = \mid + \frac{1}{2} \rangle \) (position \( z_1 \), for instance in figure 2.6(c)). As the atom moves towards the right (figure 2.6(c)), it enters a region with a different laser polarization, and the light shift decreases. The internal energy thus increases, at the expense of kinetic energy (as the atom has climbed a potential hill from \( z_1 \) to \( z_2 \) (figure 2.6(c)). Thus the atom is decelerated. At this position, there is a higher probability for
Figure 2.6: Polarisation gradient cooling for a $J = \frac{1}{2} \rightarrow \frac{3}{2}$ transition in atom. (a) Energy level diagram and Clebsch-Gordan coefficients. (b) Polarisation of a resultant laser field of two counter propagating lin-perp-lin laser field at every $\lambda/8$. (c) Sisyphus cooling picture.
the atom to be optically pumped to another magnetic substate, that is now lowest in energy \( (\varepsilon_3) \). Hereby, the internal energy of the atom is decreased again, and the excess energy is carried away by the spontaneous emission photon. This process can be repeated many times, causing the atom to decelerate, and is known by the name of Sisyphus cooling [24]. Although the absolute magnitude of the force induced by this cooling process is linear in the laser intensity, the velocity at which this maximum force is present also increases linearly with intensity. Therefore, the friction coefficient is independent of the laser intensity. As the heating by spontaneous emission is linear with the laser intensity, the temperature can be decreased indefinitely. The only limit to this cooling process shows up when the width of the velocity distribution is comparable to the recoil of one single photon [28]. This limit is called the recoil limit.

The recoil temperature is defined as the temperature at which the velocity spread of the atoms is one single photon recoil:

\[
T_{\text{rec}} = \frac{M v_{\text{rec}}^2}{k_B}
\]  

(2.11)

The significance of the recoil temperature can be understood using the wave picture for the atom. The quantum mechanical wave packet representing an atom’s motion has a de Broglie wavelength \( \lambda_{DB} \approx \hbar / M v_{\text{rec}} = \lambda \). For atoms at the recoil temperature, their deBroglie wavelength \( \lambda_{DB} \) is comparable to the optical wavelength \( \lambda_{DB} \). For \( \lambda_{DB} \approx \lambda_L \) the motion of atoms must be treated quantum mechanically. A classical treatment is valid only for atoms with temperatures much greater than \( T_{\text{rec}} \), i.e., with speeds of many \( v_{\text{rec}} \). The lower the temperature of the atom, the more important the regime of quantised motion is.

This “Sisyphus” effect is also present in another type of laser cooling, known as magnetically induced laser cooling. In this configuration, a standing wave is formed by two counterpropagating laser beams, both having \( \sigma^+ \) polarisation (figure 2.10), in the presence of a magnetic field which is perpendicular to the laser beams, providing a Zeeman shift significantly smaller than the light shift. The level scheme for an atom with \( J_g = \frac{1}{2} \rightarrow J_e = \frac{3}{2} \) is given in figure 2.7(a), and figure 2.7(b) shows the light shift potential for the same ground states. An atom initially at the position \( z_1 \), optically pumped to \( | + \frac{1}{2} \rangle \) moves to position \( z_2 \) and loses kinetic energy while climbing the potential hill. At \( z_2 \) there is a Larmor precession of the atom about the magnetic field, as the light field is zero. If the atom has precessed to the \( | - \frac{1}{2} \rangle \) state it gains less kinetic energy moving to \( z_3 \) than it has already lost, meaning that it has net energy loss. If optically pumped
Figure 2.7: MILC-Magnetically Induced Laser Cooling for a $J = \frac{1}{2} \rightarrow \frac{3}{2}$ transition in an atom. (a) Energy level diagram and Clebsch-Gordan coefficients. (b) Sisyphus cooling picture and a light-shift potential $U_p$ for the two ground state sublevels in a $\sigma^+\sigma^+$ laser field.
to \( | + \frac{1}{2} \rangle \) (point \( z_4 \)) the process starts all over again.

The other mechanism of cooling to the recoil limit is the case of counterpropagating, oppositely circularly polarised laser beams \((\sigma^+ - \sigma^-)\). In this configuration, the resulting laser field is linear at all positions, but the direction of linear polarisation rotates as a function of the relative phases of the two laser beams. Hence, a moving atom "sees" a slowly rotating electric field. In a detuned laser field, the coupling of the polarisation of the atom and the moving laser polarisation results in a population imbalance between the \( m = \frac{1}{2} \) and \( m = -\frac{1}{2} \) magnetic substates. For negative detuning and an atom moving counter propagating to the \( \sigma^+ \) polarised laser, the population in the \( m = \frac{1}{2} \) state is larger, and this means that the atom interacts more strongly with the counterpropagating \((\sigma^+)\) laser. The details on the calculation of the population differences can be found in reference [24]. Again, the friction coefficient is independent of the laser intensity.

Thus all of these mechanisms yield temperatures bounded by the recoil limit, which for helium is \( T_r = 4.08 \mu K \).

### 2.4 Velocity Selective Coherent Population Trapping (VSCPT)

One of the techniques for achieving very low temperatures is the preparation of a "dark state", which is decoupled from the laser light field. It has been shown [29] that the population of the dark state can be made velocity selective and that sub-recoil temperatures can be achieved. In this experiment, a recoil doublet was observed in the momentum distribution of atoms, for the first time. The experiment employed the \( ^2S_1 \rightarrow ^2P_1 \) transition in metastable helium at 1083 nm, 2.3 GHz to the blue of the \( ^3S_1 \rightarrow ^3P_2 \) transition we employ to cool and manipulate the atoms.

As metastable helium is one of the elements in which VSCPT is most clearly observed, we plan to use the bright beam machine (see chapter 6.6) for a detailed study of this effect, and possibly use it to precool the atoms to close to the Bose-Einstein Condensation point. Here, I will give an introduction to the mechanism.

In this scheme, the light field is formed by \( \sigma^+ \sigma^- \) laser beams, as shown in figure 2.8a. An atom, initially in the \( |g^0 \rangle \) ground state \((m = 0)\) can be excited to \( |e\pm \rangle \) and transferred to the \( |g^\pm \rangle \) ground states by spontaneous emission. Both \( |g^\pm \rangle \) states can be excited to the \( |e^0 \rangle \) state only and, since the transition \( |e^0 \rangle \rightarrow \)
2.4. VSCPT

$|g^0\rangle$ is forbidden, the atom can only make a transition to one of the $|g^\pm\rangle$ states. After some absorption-emission cycles all atoms will be optically pumped to one of these two ground states (Fig 2.8b).

When an atom in the $|g^-\rangle$ ground state absorbs a $\sigma^+$ photon and gets transferred to the $|e^0\rangle$ excited state, it absorbs momentum $+\hbar k$ as well. In the same manner, a $\sigma^-$ photon, travelling in the opposite direction, induces a transition from the $|g^+\rangle$ to the $|e^0\rangle$ state, and changes the momentum of the atom by $-\hbar k$. Thus the momentum of the atoms centre-of-mass is quantised as shown in the figure. The following three momenta form a closed family, and the average momentum of the closed family is known as the family momentum $P_f$:

$$|g^-, P_f - \hbar k\rangle$$
$$|g^+, P_f + \hbar k\rangle$$
$$|e^0, P_f\rangle.$$  

These three states are coupled to each other through the $\sigma^+\sigma^-$ laser field (absorption and stimulated emission), and through the vacuum field - by means of spontaneous emission - to other families (Fig 2.8c). A measurement of the momentum for a particular family will show two components at $P_f \pm \hbar k$, as the excited state component will have decayed before the atom reaches the detector. This separation is two times larger than is known from the saturated absorption spectroscopy measurements [30]. The VSCPT experiments [29], in which atoms are trapped in the zero momentum family, show a doublet in the measured velocity at $\pm v_{rec} = \pm \hbar k/m$ with a very small velocity spread.

This observation might be clearer if we go back to the description of the radiative velocity dependent force, as the superposition of an average damping

![Figure 2.8: The Velocity Selective Coherent Population Trapping mechanism. For details, see text.](image-url)
force and a random force (Figure 2.4). A damping force would produce a peak at \( P_f = 0 \). For a purely random force, the atom undergoes a random walk in family momentum space, changing the family momentum by an amount within \( \pm 2\hbar k \) with each step, and with the time between steps inversely proportional to the r.m.s. value of the random force.

There are states in which random force fluctuations for an atom vanish, meaning that the atom does not absorb or emit photons. These states, in which the atom is decoupled from the light field, are known as dark states for a given laser field. Typically, for the transition between the ground and excited states with the same angular momentum \( J \), the \( m_g = J \) ground state is dark in a \( \sigma^+ \) laser field, due to the fact that selection rule allows \( \Delta m = +1 \) in this case, and there is no \( m_e = J + 1 \) state to be excited to. In the same way, in linearly polarised laser light, the \( m = 0 \) ground state is dark, again because of the selection rule which allows only \( \Delta m = 0 \) for the absorption of the photon from a linearly polarised laser field. The Clebsch-Gordan coefficient \( C \) for the transition between \( m = 0 \) ground states and excited states vanishes when those states have the same angular momentum \( J \).

If we go back to the situation depicted in figure 2.8, with an atom placed in the laser field formed by two counterpropagating \( \sigma^+\sigma^- \) laser beams resonant with a \( J = 1 \rightarrow J = 1 \) transition, it is obvious that both ground states \( |g^\rightarrow\rangle|g^\rightarrow\rangle \) are coupled to the excited \( |e^0\rangle \) state in the given laser field. This, in turn, means that none of these ground states is a dark state, but it is possible to create a dark state as a superposition of the two non-dark ground states: \( |\Psi_{nc}\rangle = a|g^-\rangle + b|g^+\rangle \). Depending on the relative phase of the coefficients \( a \) and \( b \), a superposition of \( |g^\pm\rangle \) ground states can result in two orthogonal linear combinations of these states: \( |\Psi_c\rangle \), coupled to the \( |e^0, p\rangle \) with Rabi frequency \( \sqrt{2}\omega_l \) and \( |\Psi_{nc}\rangle \), the dark state. Those two orthogonal states, for any family momentum, can be expressed as:

\[
|\Psi_{nc,c}\rangle = (|g^-, P_f - \hbar k\rangle \pm |g^+, P_f + \hbar k\rangle)/\sqrt{2}.
\]  

We are not interested in any dark state, but only in one for which the family momentum \( P_f = 0 \):

\[
|\Psi_{nc}(0)\rangle = (|g^-, -\hbar k\rangle + |g^+, +\hbar k\rangle)/\sqrt{2}
\]  

which is a stationary state since the two states \( |g^\pm, \pm \hbar k\rangle \) have the same internal and kinetic energies, and since \( \langle \Psi_{nc}(0)|H_{int}|e^0\rangle = 0 \). \( |\Psi_{nc}(0)\rangle \) is an eigenstate of
the total Hamiltonian of the system. The coupling of $|\Psi_{nc}\rangle$ to $|\Psi_c\rangle$ is actually proportional to $P_f$. $|\Psi_{nc}\rangle$ is not an eigenstate of the momentum operator, so that for atoms trapped in $|\Psi_{nc}(0)\rangle$, the atomic momentum distribution presents two peaks at eigenvalues $p_{at} = \pm \hbar k$.[29] If the atom spontaneously emits a photon, and ends up in a momentum family for which family momentum $P_f \approx 0$, it is not very likely that it will again leave that family, as the time scale of change of the family momentum due to the spontaneous emission is large compared to the interaction time. As a result atoms pile up near $P_f = 0$, which is the observed narrow peak structure.

We envisage using this population trapping scheme to generate a very well collimated atomic beam for future use in atom optics experiments, and to use the coherence between the $p_m \hbar k$ momentum components for atom interferometry.

## Cooling parameters for different atoms

Cooling parameters for different atoms are given in table 2.1.

<table>
<thead>
<tr>
<th>Atom</th>
<th>$^1H$</th>
<th>$^4He^*$</th>
<th>$^4He^*$</th>
<th>$^7Li$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mass M (amu)</td>
<td>1.0078</td>
<td>4.0026</td>
<td>4.0026</td>
<td>7.016</td>
</tr>
<tr>
<td>Wavelength λ (nm)</td>
<td>121.6</td>
<td>389</td>
<td>1083</td>
<td>671</td>
</tr>
<tr>
<td>Natural width $\Gamma/2\pi$ (MHz)</td>
<td>100</td>
<td>1.7</td>
<td>1.62</td>
<td>5.8</td>
</tr>
<tr>
<td>Recoil velocity $v_{rec} = \hbar k/M$ (cm/s)</td>
<td>326</td>
<td>25.6</td>
<td>9.20</td>
<td>8.48</td>
</tr>
<tr>
<td>rms velocity at $T = T_{Doppler}$ (cm/s)</td>
<td>440</td>
<td>29</td>
<td>28.2</td>
<td>41</td>
</tr>
<tr>
<td>3D Doppler temperature (μK) $T_D = \hbar \Gamma/2k_B$</td>
<td>2390</td>
<td>40</td>
<td>39</td>
<td>140</td>
</tr>
<tr>
<td>Recoil temperature (μK)</td>
<td>1285</td>
<td>31.6</td>
<td>4.08</td>
<td>6.07</td>
</tr>
<tr>
<td>Recoil parameter $\epsilon_{rec} = \hbar k^2/2M\Gamma = \omega_{rec}/\Gamma$</td>
<td>134</td>
<td>0.2</td>
<td>26.2</td>
<td>10.8</td>
</tr>
<tr>
<td>x10$^{-3}$</td>
<td>x10$^{-3}$</td>
<td>x10$^{-3}$</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Recoil frequency/2π (kHz) $\omega_{rec}/2\pi = mv_{rec}^2/2\hbar$</td>
<td>13390</td>
<td>330</td>
<td>42.5</td>
<td>63.2</td>
</tr>
<tr>
<td>Recoil time (μs), $t_{rec} = 1/\omega_{rec}$</td>
<td>0.01189</td>
<td>0.483</td>
<td>3.745</td>
<td>2.52</td>
</tr>
<tr>
<td>Saturation intensity (mW/cm²)</td>
<td>7200</td>
<td>3.7</td>
<td>0.167</td>
<td>2.52</td>
</tr>
</tbody>
</table>

*Table 2.1: Laser cooling parameters*
## Table 2.2: Laser cooling parameters

<table>
<thead>
<tr>
<th>Atom</th>
<th>$^{20}Ne$</th>
<th>$^{23}Na$</th>
<th>$^{85}Rb$</th>
<th>$^{133}Cs$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mass M (amu)</td>
<td>19.99</td>
<td>22.99</td>
<td>84.91</td>
<td>132.9</td>
</tr>
<tr>
<td>Wavelength $\lambda$ (nm)</td>
<td>640</td>
<td>589</td>
<td>780</td>
<td>852</td>
</tr>
<tr>
<td>Natural width $\Gamma/2\pi$ (MHz)</td>
<td>8.2</td>
<td>9.7</td>
<td>6.0</td>
<td>5.2</td>
</tr>
<tr>
<td>Recoil velocity $v_{rec} = \hbar k/M$ (cm/s)</td>
<td>3.11</td>
<td>2.95</td>
<td>0.602</td>
<td>0.352</td>
</tr>
<tr>
<td>rms velocity at $T = T_{Doppler}$ (cm/s)</td>
<td>29</td>
<td>29</td>
<td>12</td>
<td>8.9</td>
</tr>
<tr>
<td>3D Doppler temperature ($\mu K$)</td>
<td>200</td>
<td>240</td>
<td>140</td>
<td>130</td>
</tr>
<tr>
<td>$T_D = \hbar \Gamma/2k_B$</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>recoil temperature ($\mu K$)</td>
<td>2.34</td>
<td>2.40</td>
<td>0.371</td>
<td>0.198</td>
</tr>
<tr>
<td>Recoil parameter $\epsilon_{rec} = \hbar k^2/2M\Gamma = \omega_{rec}/\Gamma$</td>
<td>$3.0 \times 10^{-3}$</td>
<td>$2.5 \times 10^{-3}$</td>
<td>$0.64 \times 10^{-3}$</td>
<td>$0.39 \times 10^{-3}$</td>
</tr>
<tr>
<td>Recoil frequency $\omega_{rec}/2\pi$ (kHz)</td>
<td>24.3</td>
<td>25.0</td>
<td>3.86</td>
<td>2.07</td>
</tr>
<tr>
<td>$\omega_{rec}/2\pi = mv_{rec}^2/2\hbar$</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>recoil time ($\mu s$), $t_{rec} = 1/\omega_{rec}$</td>
<td>6.54</td>
<td>6.37</td>
<td>41.2</td>
<td>77.0</td>
</tr>
<tr>
<td>Saturation intensity (mW/cm$^2$)</td>
<td>4.1</td>
<td>6.4</td>
<td>1.6</td>
<td>1.1</td>
</tr>
</tbody>
</table>
Polarisation configurations

Counter propagating linearly polarised waves:

Standing wave:
Resulting polarisation linear
Amplitude varies with position

Figure 2.9: Resultant laser field, generated as a superposition of two linear-parallel-linear counterpropagating laser beams.
Counter propagating circularly polarised waves of opposite helicity:

Standing wave:
Resulting polarisation circular
Amplitude varies with position

Figure 2.10: Resultant laser field, generated as a superposition of two counter propagating circularly polarised waves of opposite helicity ($\sigma^+\sigma^+$).
Counter propagating orthogonally linearly polarised waves:

Standing wave:
Resulting polarisation changing from linear to circular in $\lambda/8$
Amplitude constant

Figure 2.11: Resultant laser field, generated as a superposition of two linear-perpendicular-linear counterpropagating laser beams.
Counter propagating circularly polarised waves of identical helicity:

Standing wave:
Resulting polarisation linear
Polarisation direction varies with position

Figure 2.12: Resultant laser field, generated as a superposition of two counter propagating circularly polarised waves of the same helicity ($\sigma^+\sigma^-$).
Chapter 3

Transverse laser cooling of a velocity-selected sodium atomic beam

3.1 Introduction

This chapter reports on an extensive study of polarisation gradient cooling techniques. The post collimation section of the bright beam source, described later in this thesis, employs the polarisation gradient cooling mechanism to further narrow the transverse velocity spread of an already compressed, slow beam of helium atoms. The extensive study of this cooling mechanism was conducted as an experiment done in parallel with the building of our bright beam machine. The results of a detailed analysis of the experimental results gave us a better insight into the efficacy of the polarisation gradient cooling process with regard to parameters such as the cooling laser intensity, detuning and polarisation.

We employed the polarisation gradient technique to narrow the transverse distribution of a thermal sodium atomic beam. We used counterpropagating laser fields perpendicular to the atomic motion and detuned slightly below resonance. This cooling mechanism was enhanced through the use of polarisation gradients, which can lead to temperatures that are orders of magnitude lower than achievable by Doppler cooling techniques, as detailed in Chapter 2. A number of papers on this subject can be found in reference [31].

The polarisations of the counterpropagating laser fields were changed from linear-parallel-linear to linear-perpendicular-linear and circular ($\sigma^+\sigma^+$ and $\sigma^+\sigma^-$).
The transverse velocity distribution of the thermal sodium atomic beam was measured, and the results analysed, for each of these polarisations with varying laser intensities and detunings.

As the most effective cooling, or equivalently the narrowest transverse velocity distribution of the atomic beam, was observed for the case of polarisation gradients induced by two counterpropagating laser fields polarised linearly and orthogonally, we will specially consider this configuration. Such linear polarisation (fig 2.11) produces a light field whose ellipticity varies in a repeated pattern every half wavelength from linear, to circular, to orthogonal linear, to oppositely circular, back to oppositely linear. The result is that atoms which pass through the generated periodic potential experience a cooling force as a result of the Sisyphus effect [24].

The limits to polarisation gradient cooling of the transverse atomic velocity are of considerable interest both theoretically and for the applications mentioned previously. Central to this interest is the dependence of the transverse atomic temperature on the laser intensity. The early analytical treatment of polarisation gradient cooling by Dalibard and Cohen-Tannoudji [24] showed that for large intensities where diffusion effects predominate, the transverse temperature increases linearly with the laser intensity. However, the laser intensity cannot be decreased arbitrarily to produce lower temperatures as the capture range of the sub-Doppler cooling mechanism becomes smaller than a recoil. The result is that a minimum can be expected in the cooled temperature, which is predicted to be larger than the recoil temperature $T_r$, defined by

$$k_B T_r = \frac{k^2 k^2}{2M}$$

(3.1)

where $k = 2\pi / \lambda$ ($\lambda$ is the wavelength), $M$ is the atomic mass, and $k_B$ and $\hbar$ are the Boltzmann and Planck constants respectively.

A quantum treatment of this problem in the low excitation limit by Castin et al. [32] confirmed the linear dependence of the transverse temperature at higher intensities. Instead of a temperature, they used the average transverse kinetic energy, which is related to the mean momentum, $p_{r.m.s.}$ by

$$E_K = \frac{p_{r.m.s.}^2}{2M}$$

(3.2)

The intensity $I$ is related to the light shift potential $U_0$ in the two beams (nor-
malised by the recoil energy $E_r = h\nu_r$) by

$$\frac{U_0}{h\nu_r} = f \frac{\Delta}{\nu_r} \frac{s}{1 + 4\Delta^2/\Gamma^2}$$ \hspace{1cm} (3.3)

where $\nu_r$ is the recoil frequency (24.9 kHz), $\Delta$ is the interaction laser detuning, $\Gamma$ is the natural linewidth of the transition and $s$ is the single beam saturation parameter $I/I_0$ ($I_0 = 6\text{mW/cm}^2$ is the saturation intensity for the sodium transition). The factor $f$ is given by [32, 33]

$$f = \frac{(2F+1)(F+1) - 1}{(2F+1)(F+1)} = \frac{14}{15} \text{ for the } F = 2 \rightarrow F = 3 \text{ transition in sodium}$$ \hspace{1cm} (3.4)

Eqn. 3.3 is valid in the low excitation limit ($\Delta \gg \Gamma$).

At high intensities ($U_0$ larger than the value at the minimum kinetic energy) and for the one-dimensional cooling case for sodium considered here, Castin et al. [32] predict a slope of 0.14 for the transverse kinetic energy as a function of the light potential $U_0$. This slope is to be compared with a value of 0.19 from the simple analytical treatment [24].

The minimum in the kinetic energy is again predicted, yet it appears to occur at a normalised potential that is relatively independent of the atomic species, at around $U_0 \sim 100E_r$ (for sodium, at $U_0 \sim 105E_r$). Below this value, the kinetic energy increases rapidly with decreasing intensity. The minimum kinetic energy was shown to be relatively independent of laser detuning for $\Delta \gg \Gamma$, and for sodium yields a value of $E_K \sim 40E_r$.

The question arises as to whether the mean kinetic energy can be used to assign a temperature to the cooled atoms. As Castin et al. point out, if the momentum distribution is a Gaussian then a temperature can be assigned, since in this case $p_{\text{r.m.s.}}/M$ is equal to the half width of the velocity distribution at the $1/\sqrt{\epsilon}$ point. However, both in their quantum Monte Carlo calculations and in their analytical calculations, these quantities diverge, with the result that the minimum in the $1/\sqrt{\epsilon}$ width of the velocity distribution (as opposed to $p_{\text{r.m.s.}}$ curve) occurs at a value for $U_0 \sim 25E_r$ in the case of cesium. The minimum velocity width at this value of $U_0$ for a detuning of $\Delta = -15\Gamma$ is $v_{\text{min}} \sim 3\nu_r$, i.e an energy of $\sim 9E_r$.

A number of experiments using trapped atoms with long interaction times and at large detunings have verified the theoretical predictions of these models,
confirming, for the case of three dimensions, a linear dependence of the temperature on the light potential [34, 35, 36]. Experiments on one dimensional cooling of trapped atoms have also demonstrated the expected linear dependence on the light potential [37, 38].

However, the situation is less conclusive for the one-dimensional transverse cooling of a thermal atomic beam. The width of the transverse velocity distribution has not shown the expected dependence on the light potential in experiments using a rubidium atomic beam [39]. This experiment revealed a lower slope for $E_K$ versus $U_0$ than expected from the previously mentioned theoretical predictions: indeed an almost flat behaviour for the transverse kinetic energy at higher intensities was observed. These claims have been revised in some details recently for the case of one dimensional cooling of chromium thermal atomic beam [40]. Their measurements have shown the general trends expected from low excitation approaches, which predicted improved cooling and reduced kinetic energies along the laser axes as the light shift is increased from zero, reaching a characteristic minimum energy spread before rising due to diffusion heating effects at higher laser intensities.

However, as was pointed out by Bergeman [33] and later by Doery et al. [41], there are a number of complicating factors in beam experiments which inhibit the interpretation of such results. Firstly, in contrast to atom trap studies, the interaction time for atomic beam experiments is usually limited by the maximum available laser power, which determines the maximum interaction length. Typical atomic beam interaction times are less than a few thousand natural lifetimes. In such situations the transverse velocity distribution seldom reaches the steady state limit (particularly where large transverse velocities are present). Consequently, careful quantum Monte Carlo modelling is required to determine the expected final spatial distribution in the atomic beam [33, 41].

The difficulties of including other non-ideal experimental effects such as the shape of the initial velocity distribution, the laser intensity distribution and the interaction transit time, must be addressed. The result is that a decisive comparison between experiment and theory is very difficult. For instance experimental results for He* [42] show that in the high intensity regime, velocity distributions are typically non-Gaussian in shape. There are two velocity regimes: one associated with the steep force function near $v = 0$ from Sisyphus cooling and one with the less steep Doppler force which extends to higher velocities. Doppler cooling can greatly increase the velocity capture range, and become more significant for
3.1. **INTRODUCTION**

Light atoms such as He*. The optimum detuning in this experiment, determined in such a way that it maximizes the capture velocity while avoiding a gap in $F(v)$ between the region of the large Doppler force and the Sisyphus region, as in the previously mentioned chromium experiment [40] is found to be approximately $\delta = -\Gamma$.

The work reported here provides data for the transverse cooling of a sodium atomic beam, using counterpropagating laser beams with lin⊥lin polarisation. Various other laser polarisation configurations were also studied, but orthogonal linear polarisations were found to produce the most effective cooling.

A distinguishing feature of the present experiments is the use of a longitudinal velocity-selective detection system which enables a direct determination of the transverse velocity distribution. In previous experiments [40] the full thermal velocity distribution is used. This complicates the interpretation of the experimental results, not only atoms with a different longitudinal velocity have different interaction times with the laser, but also have a different angle with the beam axis for the same transverse velocity. The technique used in the present experiment simplifies interpretation of the experimental results by eliminating these problems.

In the present experiments, the detection region was illuminated using a second laser incident at an angle to the atomic beam in order to select a single longitudinal velocity group. The spatial profile of the resulting laser-induced fluorescence then yielded a direct measure of the transverse velocity profile, with sub-recoil resolution. By placing a knife edge immediately following the interaction region, the divergence of the atomic beam downstream from the knife edge was used to determine the transverse velocity distribution of the cooled atoms. The transverse temperature was therefore derived directly due to the separation of the longitudinal and transverse velocities.

The results reported here for the case of one dimensional cooling of sodium are again at variance with current polarisation gradient theory. Indeed, despite the existence of other non-idealised conditions that include finite interaction times, nonuniform laser field intensity distributions, and optical pumping processes in a multilevel system, we managed to achieve transverse temperatures more than an order of magnitude less than the minimum predicted by infinite interaction time models [24, 32]. Once more the experiments reveal a very weak dependence of the transverse temperature on the laser intensity.
CHAPTER 3. TRANSVERSE COOLING...

3.2 Experiment

The experimental configuration was similar to that used in previous experiments in this laboratory [43, 44] and is shown in fig 3.1. A sodium atomic beam was produced by a thermal oven source operating at a temperature of 400–450°C and with an exit aperture of 1 mm diameter. After passing through a skimmer, the beam was apertured by a 1 mm (height) × 3 mm (width) slit located 350 mm downstream from the oven, that also served as the entrance port for the differentially pumped interaction chamber. The 1 mm atomic beam height restriction was used to minimise the variation of the interaction laser intensity across the atomic beam in the vertical direction. The 3 mm width (combined with the 1 mm oven aperture) provided a broad range of transverse velocities for the cooling experiment. The interaction chamber itself was located within a full set of Helmholtz coils to minimise the effect of stray magnetic fields (residual magnetic field B < 1 μT). Light from a stabilised ring dye laser (the interaction laser - a Spectra physics 380D operated with R6G) formed the transverse cooling beam. The beam passed
3.2. EXPERIMENT

![Level structure of sodium](image)

*Figure 3.2: Part of the level structure of sodium. The two first order sidebands generated by the EOM excite the $F = 2 \rightarrow F = 3$ and $F = 1 \rightarrow F = 2$ transitions respectively, as indicated by the arrows.*

through an 856 MHz electro-optic modulator and the two first-order sidebands were used to generate the cooling and repumping laser frequencies, as indicated in figure 3.2. Each of these first order sidebands contained $\approx 1/3$ of the total laser power. An anamorphic prism pair was used to magnify the horizontal dimension of the laser beam by a ratio of 4:1. A telescope system was employed to further increase the dimensions of the interaction laser beam, and was also used to collimate the laser to $< 0.5$ mrad. The intensity in the interaction region could be varied by rotation of a half wave plate placed before a linear polariser. The maximum laser intensity averaged over the interaction region was 40 mW/cm². However, only 1/3 this total intensity is at resonance with the $F = 2 \rightarrow F = 3$ cooling transition. This corresponds to a maximum, single beam, on-resonance saturation parameter of $s \sim 2$ at the centre of the beam.

An aperture with a 15mm horizontal width was placed in the laser path prior to the interaction region at a distance of 75 mm from the atomic beam (outside the interaction chamber). The edges of the aperture defined the length of the
interaction region. The collimated laser beam intensity profile was measured very precisely using a 25 μm pinhole and a power meter mounted on a micrometer-driven translation stage on the far side of the interaction region. The pinhole was scanned horizontally across the laser beam, and the resulting profile is shown in the inset to fig.3.1.

The structure in the measured laser intensity profile results from diffraction and interference effects. This structure may also be present in similar experiments elsewhere, but may not be noticeable unless measured using such a high resolution, two dimensional technique. The translating slit method often employed averages out variations in one dimension and may smooth out such spatial structure.

In the present experiment, the structure could not be eliminated easily without the use of spatial filtering, which would have reduced the laser intensity to impractical levels. The period of the structure however, is very small compared to an atomic oscillation period in an optical potential well, and the structure is therefore believed to have only minor implications for the cooling results (see Discussion). The rising/falling edge of the light field had a width of 300 μm, which was made as small as possible to minimise the effect of adiabatic cooling [38].

Upon passing through the interaction region the laser was retro-reflected from a rotatable mirror located outside the interaction chamber 100 mm from the atomic beam. Depending upon the cooling method used, polarising elements could be placed between the mirror and the exit window to define the interaction field. The polarisation schemes used were \( \sigma^+ - \sigma^- \), \( \sigma^+ - \sigma^+ \), linear and orthogonal linear (lin_Llin - the configuration shown in fig.3.1).

The laser beam entering the interaction region was aligned horizontally to within a few mrad. Perpendicular alignment to the atomic beam axis is achieved by using a corner cube reflector on the far side of the interaction region to produce a second fluorescence spot in the atomic beam. By detuning the laser frequency appropriately, the relative brightness and spatial distribution of the fluorescence spots could be monitored, a symmetric behaviour indicating that perpendicular alignment of the incoming laser had been achieved. The perpendicularity of the laser beam reflected from the rotatable mirror was achieved by ensuring that the cooled atomic beam profile was spatially symmetric with respect to the centroid of the uncooled beam.

The interaction laser was tuned to the sodium \( \{3s\}^2S_{1/2} \ (F=2) \rightarrow \{3p\}^2P_{3/2} \)}
3.2. EXPERIMENT

(F=3) transition at 589.1 nm. The laser frequency was referenced to the signal from a sodium absorption cell. The zero detuning of the laser could also be determined to within ±1 MHz by the rapid change from cooling to heating effects as the laser goes through resonance. A high finesse spectrum analyser with a 1.5 GHz free spectral range and a stability of ±2 MHz over the observation period was then used to measure the interaction laser detuning with a repeatability of < Φ/4.

The spatial profile of the atomic beam was measured by a velocity selective laser fluorescence system located in a separate vacuum chamber 1.64 m from the interaction region. The system consisted of a second, independently tunable laser (a single mode Coherent 699-21 operated with R6G) which was tuned to the transition between the \{3s\}^2 S_{1/2} (F=2) hyperfine ground state and an excited state. The laser induced fluorescence (LIF) produced as the atoms decayed into both the \{3s\}^2 S_{1/2} (F=2) and \{3s\}^2 S_{1/2} (F=1) states yielded a signal proportional to the density of atoms in the detection region in the F=2 state and resonant with the detection laser. (Strictly speaking, there are three accessible upper hyperfine levels, but the fluorescence signal is dominated by the F=2 → 3 transition.) This state selective detection system is therefore used to detect only those atoms which interacted with the cooling laser field.

In order to provide velocity selectivity, the detection laser crosses the atomic beam at 74° to the propagation axis of the atomic beam. For sufficiently low detection laser powers (below saturation), the longitudinal velocity resolution is determined by the atomic linewidth (9.7 MHz) which was much greater than the laser bandwidth (~1 MHz). The combination of intersection angle and detection bandwidth yields a longitudinal velocity resolution of ±10 m/s. The detection laser was detuned by 320 MHz to select a longitudinal velocity of 700 m/s. This was just below the peak of the Maxwellian velocity distribution in order to detect slower atoms with a longer interaction time, while at the same time providing reasonable signal levels. The detection laser detuning was verified before and after each experiment by observing the shift of the fringes on a scanning Fabry-Pérot interferometer when tuning away from resonance.

The LIF was detected by a linear, 1024 element Reticon photodiode array via a dual-lens achromat imaging system with unity magnification. The detection laser was focused in the horizontal direction using a cylindrical lens to provide a fluorescence region with a small depth of field (and hence sharper image) for the detection system.
The size of the detector pixels (25 μm) determines the spatial resolution of the apparatus. The Reticon was parallel to the detection laser beam i.e. at 74° to the atomic beam axis, and the 1024 element linear array was aligned horizontally. The distance to the interaction region yields an angular resolution for each pixel of \(\approx 15 \mu\text{rad}\). This corresponds to a transverse velocity resolution of 1 cm/s (\(\approx \frac{1}{3}\hbar k\)) for a longitudinal atomic velocity of 700 m/s.

The position of the detector relative to the fluorescence region can be adjusted using a two axis positioning system to focus and translate the detector in the horizontal direction. In order to optimise the focus, a slit \(\approx 15 \mu\text{m}\) wide and \(\approx 1 \text{ mm}\) high can be moved into the atomic beam immediately upstream of the detection region. The image of the resulting 15 μm wide fluorescence region can then be optimised using the focusing mechanism, yielding a FWHM of 3 pixels. Given that the fluorescence region is itself \(\approx 1 \text{ pixel}\) wide, this measurement determines the instrument resolution of the entire detection system due to all sources of systematic error to be than \(\approx 1 \text{ pixel r.m.s.} (\pm\frac{1}{3}\hbar k)\).

The detector array accumulates signal from the detection region and the output is sampled after an integration time of 2 s. A sequence of 5 traces is stored and the result averaged to improve the signal-to-noise ratio. Background subtraction techniques are further employed to reduce systematic contributions to the noise level. Detailed spatial profiles of the transversely cooled atomic beam in the far field were then obtained using this system.

In general, the spatial distribution consisted of a well-defined peak of transversely cooled atoms on a broad background of atoms which had experienced significantly less interaction with the cooling laser field. To determine the transverse velocity distribution of the cooled atoms, we used a knife edge to partially block the atomic beam 160 mm downstream from the laser cooling region [40]. After the knife edge, the atoms traversed 1.48 m until they were illuminated by the detection laser.

The knife edge was aligned vertically using diffraction from a helium neon laser to produce a diffraction pattern that was horizontal to better than a few mrad, which was all that was needed to ensure parallelism of the knife edge with the 25 × 1000 μm vertical elements of the photodiode array.

The knife edge provides the effect of a point source. To determine the efficacy of the laser cooling process, the one dimensional angular distribution of the cooled beam as \(f(\theta)\), where \(\theta\) is the angle to the atomic beam axis, in the direction of the laser beams is of interest. This can now be derived from the measured spatial
3.2. EXPERIMENT

Figure 3.3: (a) Atomic beam spatial profile using $\text{lin}_\perp\text{lin}$ polarisation, with $\Delta = -0.6\Gamma$ and a single beam intensity of 13 mW/cm$^2$. Also shown is the fitted sum of the three error functions, each individually shown in (b). For the fit (solid line in (a)), only the left hand part of the data was used.

Figure 3.4: The derivative of the fitted curve (line) and of the raw data.
profile as follows: The measured spatial fluorescence intensity $I$ can be written as:

$$I(x) = \int_{\theta_{\text{min}}(x)}^{\pi/2} f(\theta) d\theta$$  \hspace{1cm} (3.5)$$

where $I(x)$ is the fluorescence intensity at distance $x$ from the beam axis and $\theta_{\text{min}} = \tan^{-1}(x/1.48)$, assuming the knife edge blocks the atomic beam for $x < 0$. Taking the derivative of this profile then provides the angular distribution $f(\theta)$.

Because we have selected a single longitudinal velocity group, then the divergence angle is linearly related to the transverse velocity. That is, an atom following a path at an angle $\theta$ to the beam axis at longitudinal velocity $v_l$ must have transverse velocity $v_t \theta$. Hence, from $f(\theta)$ we have directly the transverse velocity distribution, because $f(v_l) = f(\theta)$. The width of the derivative of $I(x)$ (i.e. $f(\theta)$) is a measure of the width of the transverse velocity distribution.

The knife edge was placed at the peak of the cooled atomic spatial distribution. By differentiating the resulting spatial profile for atoms detected in a specific longitudinal velocity group (defined by the detection laser frequency), the transverse velocity distribution of the cooled atoms in that group could be measured directly.

The knife edge technique has distinct advantages over the, more conventional, narrow slit technique. An important advantage in a sodium experiment is that a narrow slit can easily get clogged by the sodium. This can be overcome by heating the slit material, but this may cause other problems, such as warping of the slit material. Another advantage is that the resolution is determined only by the detection system, not by the slit width.

A disadvantage of the knife edge method is the added complexity of the analysis. The structure on the right side of figure 3.3 has to be ignored as it is not caused by the knife edge, but by the finite capture range of the cooling process. Hence the maximum of the central feature in this figure was taken as the cut-off point for the data analysis.

### 3.3 Results

The spatial distributions of atoms near the peak of the longitudinal velocity distribution ($v_l = 700 \text{ m/s}$) were measured as a function of laser polarisation, intensity and detuning (see figures 1–10 in the Appendix). The most efficient cooling po-
3.3. RESULTS

lolarisation was found to be linear, and it is the results of these measurements that are presented here.

An example of the spatial profile measured in the experiments is shown in fig. 3.3a. The profile consists of a sharp peak of cooler atoms on a broad background, and the peak is cut by the knife edge with the penumbral region falling off to the left. In order to minimise the effect of noise on the differentiation of the spatial profile, the penumbral region for the cold peak is fitted by a combination of error functions (for which the derivative is a Gaussian) using a least squares fitting routine.

The first iteration used two error functions, one to simulate the cooler beam with the other representing the broad, hotter background. However, in order to achieve a good fit, it was found necessary to introduce a further error function. The three individual error functions are shown in fig. 3.3b and their resultant fit is shown in fig. 3.3a.

The derivative for the fitted curve is shown in fig. 3.4 together with the derivative of the raw data. As can be seen, the main Gaussian fits the narrow central peak reasonably well, while the small secondary peak fits the shoulder on the left hand side. The broad Gaussian corresponding to the hotter atoms is not evident as it is much flatter on this scale.

The results of the calculations by Castin et al. [32] indicate that such velocity distributions may not always be Gaussian, and hence cannot always be characterised by a temperature. Depending upon whether a compact (small p\text{r.m.s.}) momentum distribution or a narrow central peak (small $1/\sqrt{\varepsilon}$ half width) is used as a fitting criteria, both the optimal value of $U_0$ and the minimum energy/temperature may vary by a factor of 4 as mentioned previously.

However, as can be seen from fig. 3.4, a Gaussian fit to the narrow central peak is a good approximation. The $1/\sqrt{\varepsilon}$ half width of the spatial profile for the cooled atoms is then converted to transverse velocity half widths $v_t$ directly using the pixel calibration from the previous section. The temperature $E_K = \frac{1}{2}Mv_t^2$ is then plotted as a function of the light potential, and is shown in fig. 3.5. Also shown is the area under the derivative curve (proportional to the total number of cooled atoms) as a function of the light potential.

The laser intensity is expressed in normalised units for the light shift potential divided by the recoil energy. However, at the small detunings employed in the present experiment ($\Delta \sim \Gamma$) the previous expression for the light shift $U_0$ is not valid. A more correct expression for the normalised light shift potential in this
CHAPTER 3. TRANSVERSE COOLING...

Figure 3.5: Temperature of cooled atoms (solid squares) normalised to the photon recoil energy ($E_r$) and area under cooled peak (circles) as a function of $U_p$ ($\Delta = -0.6\Gamma$). The solid line represents a best linear fit to the temperature points for $U_p > 70E_r$. Note that the Doppler limit $E_D = 100E_r$ for sodium.

parameter range has been derived for the case of a two level atom by Gordon and Ashkin [45]:

$$\frac{U_p}{\hbar \nu_r} = f \frac{\Delta}{\nu_r} \ln \left\{ 1 + s/(1 + 4\Delta^2/\Gamma^2) \right\}$$  \hspace{1cm} (3.6)

Note that the expression in equ. 3.6 is twice the single laser potential, as is also the case for $U_0$ in equ. 3.3. However, the logarithmic term in equ. 3.6 takes into account the excited state character of the atomic wavefunction near resonance. This is usually neglected at large detunings by omitting "ln" from the expression, as is done by Castin et al. [32] in the weak excitation limit ($\Delta \gg \Gamma$). By contrast, Equ. 3.6 gives a better approximation at small detunings ($\Delta \geq \Gamma$) and in the adiabatic limit.

The detuning dependence of the normalised temperature is shown in fig. 3.6. Here, a sequence of detunings is performed at three different laser saturation values $s$, and the temperature plotted as a function of $U_p(s, \Delta)$. As was the case for fig. 3.5, the vertical error bars are obtained from the least squares fitting routine.
3.4 Discussion

The results shown in fig. 3.5 confirm the measurements obtained elsewhere for chromium [40] and rubidium [39], namely that the cooling efficiency (number of cooled atoms) increases monotonically with intensity, and that the transverse temperature is remarkably insensitive to laser intensity. Over the range of normalised potentials from $U_p = 70 - 140E_r$, there is only a slight increase in temperature. At low values of the potential ($U_p < 70E_r$) the temperature rises sharply with decreasing potential. In this regard, the results are in qualitative agreement with theory.

As indicated by the discussion in ref. [32], the minimum turning point occurs at a value of the potential for which $U_0 = 105E_r$ or $25E_r$, depending upon whether the velocity distributions can be represented by a single Gaussian or otherwise. Here, we fit the velocity distributions (fig. 3.3c) to a sum of three Gaussians. We therefore need to compare the temperature (as derived from the half width at the $1/\sqrt{e}$ point, of the cold part of the measured velocity distribution) to the latter theoretical value for which the minimum turning point is $25E_r$.

The quantitative comparison is not very good. The slope of the straight line fitted to the experimental curve for $U_p > 70E_r$ is 0.007, much smaller than the theoretical value of 0.14. The experimental turning point value for $U_p$ is $\sim 70E_r$.
c.f. 25 $E_r$. The discrepancy may be caused by the fact that the theoretical curve was calculated for large detunings ($\Delta = -15\Gamma$), whereas the measurements in figure 3.5 were obtained at small detunings ($\Delta = -0.6\Gamma$).

The effect of different combinations of laser intensity and detuning leading to the same value for $U_p$ is illustrated in fig. 3.6. At large detunings ($\Delta \gg \Gamma$), the extent of cooling should simply scale on a single curve with $U_p$. However, at the small detunings used here, fig. 3.6 indicates that this is not the case. At varying laser intensities, the dependence of $E_K$ on $U_p$ is different. The expected minimum in temperature occurs at $U_p \approx 50E_r$, except at the lowest intensity, for which a minimum was not measured. The slope changes dramatically, from 0.09 for $s = 0.9$ to 0.02 for $s = 2.2$. Consequently, it is difficult to compare the theoretical results, obtained at large detuning, and plotted against $U_0$, to the present data, plotted against $U_p$.

The most remarkable result, however, is that the minimum temperatures obtained are much cooler than predicted by earlier theories [24, 32]. The lowest temperature observed here is $\sim 1.8E_r$, which is more than an order of magnitude less than the minimum theoretical temperature as determined from p.r.m.s. ($40E_r$), and a factor of $\sim 5$ less than the narrow central peak criterion [32]. The lowest velocity widths achieved (4 cm/s) are approaching the recoil velocity (2.95 cm/s), which is as low as has been achieved by, for example, adiabatic cooling [46].

The question then arises as to whether adiabatic cooling can contribute to the low temperatures observed in the present experiment. The result of a simple estimate shows this to be unlikely.

The potential indicated by the laser intensity profile shown in the inset to figure 3.1 decreases at the edge of the aperture over a distance of $\sim 300\mu m$, which for 700 m/s atoms gives a transit time $t_t \sim 400$ ns.

To satisfy adiabaticity one must have $|\omega|/\omega = \epsilon \omega$ [47], where $\epsilon \ll 1$ and $\omega$ is the oscillation frequency of atoms in the potential wells, given by

$$\omega = \frac{E_r}{\hbar} \sqrt{\frac{4\hbar|\Delta||s|}{E_r(1 + 4\Delta^2/\Gamma^2)}}$$

(3.7)

In our case, we have an upper limit for $\omega \sim 4 \cdot 10^6$, leading to $\epsilon > 0.3$. The values for $\epsilon$ considered by Kastberg et al. [47] to be necessary for adiabatic cooling are in the range $0.02 < \epsilon < 0.2$. The upper limit for $\omega$ above corresponds to at most one full oscillation period during the transit time through the falling edge.

Consequently, we believe that in the present experiment there is insufficient
time for the atoms to experience adiabatic cooling. Similar arguments can be made to discount adiabatic cooling in the structure evident on the laser intensity profile, for which the length scales are even smaller. While a complete Monte Carlo calculation of the atomic trajectories is required to fully account for the effects of spatial inhomogeneity, the indications are that adiabatic cooling is not responsible for the low temperatures observed.

Comparison with the full quantum Monte Carlo treatment of Bergeman [33] is difficult, since the case of sodium was not considered. However, similar calculations for rubidium for similar interaction times \((1300\Gamma^{-1})\) yield temperatures approaching the lowest values measured here in sodium. A direct comparison of theoretical calculations with the measured velocity distributions under the current experimental conditions is therefore needed.

Finally, it remains to consider the existence of the third but much smaller Gaussian required to fit the experimental velocity distributions. This suggests the presence of a third (transverse) velocity group of atoms, slightly displaced from the central peak of cooled atoms, and also significantly cooled (with widths of a few \(v_r\)). This shoulder on the transverse velocity distribution may arise from a slight imbalance in the laser intensity, introduced by the optical elements through which the laser beam passes before being retroreflected. For the \(\sigma^+\sigma^-\) polarization configuration these effects have been discussed extensively by Werner et al. [48], and yield similar shoulders in the velocity distribution.

3.5 Conclusion

Measurements have been performed with sub-recoil resolution on the transverse cooling of a longitudinally velocity-selected sodium atomic beam in a lin.Llin laser field. A minimum temperature occurs as a function of \(U_p\), in qualitative agreement with theoretical predictions.

However, at the small detunings used, the expected scaling with \(U_p\) appears to be complex, which makes quantitative comparison with theory difficult. At a fixed, small detuning, the temperature nevertheless appears to be less sensitive than expected for values of \(U_p\) larger than the value at the minimum temperature.

Furthermore, the lowest kinetic energies obtained \((E_K \sim 1.8E_r\text{ i.e. } v_t \sim 1.4v_r)\) are significantly lower than expected. It is unlikely that adiabatic cooling effects are responsible for such low temperatures.
However, quantum Monte Carlo calculations for other atomic species do indicate similar temperatures for polarisation gradient cooling. These calculations also suggest that careful account needs to be taken of the effects of variations in the spatial intensity and velocity distributions in the calculation of the cooled atomic distributions [33, 41].

Finally, there is some experimental evidence of a second cooled transverse velocity group, possibly the result of imbalance in the counterpropagating laser beams.

In order to make a direct comparison with all features of the present data, a full quantum mechanical Monte Carlo calculation of the atomic trajectories using a polarisation gradient model for sodium under the same experimental conditions is required.

The experiment could improved, for example by using an in-vacuum mirror to remove the imbalance in the intensity of the two counter-propagating beams, by improving the homogeneity of the two beam profiles, and by using in-vacuum knife edges to sharpen the transition from dark to light.

Appendix Chapter 3

In this appendix we will present data showing different cooling efficiencies for different polarisation schemes of the cooling laser. We begin with the graphical presentation of data for a lin perp lin polarisation of the cooling laser followed by the data for other polarisations for which the cooling was found to be less efficient.
3.5. CONCLUSION

Figure 3.7: Transverse spatial distribution of Sodium atoms at a longitudinal velocity of 591 ms\(^{-1}\) in the lin\perp lin\ laser field, with laser power of 102 mW, for a series of detunings ranging from \(-0.2\Gamma \text{ to } 2.5\Gamma\)

Figure 3.8: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of 591 ms\(^{-1}\) in the lin\perp lin\ laser field, with laser power of 75 mW, for a series of detunings ranging from \(-0.2\Gamma \text{ to } 2.5\Gamma\)
Figure 3.9: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of $591\text{ms}^{-1}$ in the lin x lin laser field, with laser power of $50.3\text{mW}$, for a series of detunings ranging from $-0.2\Gamma$ to $2.5\Gamma$.

Figure 3.10: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of $591\text{ms}^{-1}$ in the lin parallel lin laser field, with laser power of $102\text{mW}$, for a series of detunings ranging from $-0.2\Gamma$ to $2.5\Gamma$. 
3.5. CONCLUSION

Figure 3.11: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of $591 \text{ms}^{-1}$ in the linear laser field, with laser power of 75 mW, for a series of detunings ranging from $-0.2\Gamma$ – $2.5\Gamma$.

Figure 3.12: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of $591 \text{ms}^{-1}$ in the linear laser field, with laser power of 50 mW, for a series of detunings ranging from $-0.2\Gamma$ – $2.5\Gamma$. 
Figure 3.13: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of 591ms$^{-1}$ in the $\sigma^+\sigma^-$ laser field, with laser power of 105 mW, for a series of detunings ranging from $-0.2\Gamma - 2.5\Gamma$.

Figure 3.14: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of 591ms$^{-1}$ in the $\sigma^+\sigma^-$ laser field, with laser power of 105 mW, for a series of detunings ranging from $-0.2\Gamma - 2.5\Gamma$. 
3.5. CONCLUSION

Figure 3.15: Transverse spatial distribution width of Sodium atoms at a longitudinal velocity of 591ms⁻¹ in the $\sigma^+\sigma^-$ laser field, with laser power of 50 mW, for a series of detunings ranging from $-0.2\Gamma - 2.5\Gamma$.

Figure 3.16: Transverse spatial distribution width of Sodium atoms is found to be 3.17 recoil for atoms at velocity of 591ms⁻¹ in the lin prp lin laser field, with laser detuning of 260 MHz, for a series of laser powers ranging from 33-84 mW.
Chapter 4

The bright beam apparatus

4.1 Introduction

This chapter reports on the construction of a bright atomic beam of metastable helium atoms using a combination of laser cooling techniques. The atomic beam, generated by a liquid nitrogen cooled metastable atom source, is collimated, longitudinally slowed and spatially compressed by Doppler cooling forces. A schematics of the apparatus is shown in figure 4.1. Diode lasers are used to excite the $^2S_1 \rightarrow ^2P_2$ transition at 1083 nm for all three laser cooling stages. This chapter
4.2. LN$_2$ COOLED METASTABLE ATOM SOURCE

describes all stages of the bright beam machine and is organised as follows:

- The liquid nitrogen cooled source of metastable helium atoms is described in section 4.2.

- In section 4.3 the initial collimation of the atomic beam using effectively curved wavefronts to capture a large fraction of the atoms emerging from the source, is described.

- The Zeeman slower, in which the collimated atomic beam is slowed from $\sim$900 ms$^{-1}$ to $\sim$100 ms$^{-1}$ is described in section 4.4.

- Finally, in section 4.5 the final compression stage, in which the atomic beam is compressed to a small diameter to produce a well collimated beam, is described and the overall performance is compared with the design aims.

4.2 LN$_2$ Cooled metastable atom source

As we shall see later, a first requirement for a bright beam source is to cool the metastable atom source to liquid nitrogen (LN$_2$) temperatures. We use a conventional DC discharge source with the discharge running through an expansion nozzle into the source vacuum chamber. The design is based on a similar source for metastable neon atoms, which was developed at the University of Eindhoven [49]. As the gas flow rate in the source is small, and the nozzle is heated by the discharge current running through it, the most important design consideration to achieve a low temperature atomic beam is effective cooling of the expansion nozzle. This gives rise to two contradicting requirements for the nozzle plate:

1. it needs to be constructed from an electrical insulator and
2. it needs to have good heat conduction.

These requirements are fulfilled by the ceramic boron-nitride. Cooling is achieved with a co-axial, in-vacuum, stainless steel LN$_2$ reservoir, which surrounds the helium inlet line. The LN$_2$ reservoir is gravity fed from a large external dewar.

The cathode of the discharge is contained in a glass tube that can be pumped by an additional rotary pump to remove the hot gas around the cathode [50]. The cathode consists of a tungsten needle, as shown in figure 4.2. Its surface area can be increased by adding a nickel foil around the needle. During standard
operation, the cathode is held at -400 V with respect to ground. Breakdown to
the metal LN$_2$ reservoir, which is necessarily at ground potential, is prevented by
a Macor sheath in which a long, spiral channel has also been machined for the gas
inlet. The anode is formed by a stainless steel ring, 5 mm from the nozzle, and is
kept at +400 V. The discharge current is about 2 mA. To start the discharge, a
-5 kV start pulse is added to the cathode voltage, and the source is allowed to run
at a higher current (10 mA) for a few minutes. Then, to reduce the temperature,
the current is reduced to its operational value. The source pressure in operating
conditions is 35 Torr, and the nozzle has a diameter of 300 $\mu$m. This combination
has been found to yield the lowest temperatures and axial velocities.

The source chamber is pumped with a 2000 ls$^{-1}$ diffusion pump with a LN$_2$
cooled baffle (Varian type 184). The background helium pressure during opera­
tion is $\sim 10^{-5}$ Torr. The metastable beam passes through a skimmer into the
differentially-pumped collimation chamber.

The metastable atom beam flux obtained in typical conditions is $\sim 10^{14}$ sr$^{-1}$s$^{-1}$,
at a temperature of 80 K. The average beam velocity is then $\sim 900$ ms$^{-1}$, with
a spread of about $\sim 200$ ms$^{-1}$. However, the source can yield as much as $10^{15}$
Figure 4.3: The source temperature as a function of the driving pressure and the discharge current

sr^{-1} s^{-1} when run at a higher (15 mA) discharge current and a higher pressure (100 Torr).

In figure 4.3 the characteristics of the source are shown as a function of both the discharge current and the helium inlet pressure. The data points are obtained from a fit to the velocity distribution as measured in a time-of-flight experiment. For this experiment, the atomic beam is chopped with an in-vacuum mechanical chopper, with a slit width of 1 mm. The chopper wheel has a diameter of 100 mm and is driven by an inexpensive DC motor. It spins at about 50 revolutions/second. The atoms are detected using a channeltron electron multiplier, 40 cm downstream of the chopper, and the time-of-flight distribution is recorded on a TRACOR TN7200 multi-channel analyser. The spectrum obtained contains two signal peaks: an instantaneous peak due to UV photons from the discharge and the atom peak, which occurs at a later time. The former is used to determine the zero of the time scale \( t_0 \). A simple transformation \( v = L_{TOF}/(t - t_0) \) then yields the velocity distribution, where \( L_{TOF} \) is the flight length.

The velocity distribution is then least-squares fitted to a standard, displaced Gaussian, with an average velocity \( \langle v \rangle \) and a spread \( u \), as is usually obtained from a supersonic beam source [51]. The source temperature can then be extracted from [51]:

\[
\frac{5}{2} k_B T_{\text{source}} = \frac{3}{4} m u^2 + \frac{1}{2} m \langle v \rangle^2
\]  

(4.1)
At low pressure (30 Torr) and low discharge current (2 mA) the source temperature approaches the liquid nitrogen temperature. The average velocity \( \langle v \rangle \) is then about 900 m s\(^{-1}\). No further improvement was found by pumping on the cathode region of the source to remove the hotter gas around the cathode. Consequently, this feature is not used at present.

### 4.3 Collimation

We use Doppler cooling to collimate the initially diverging atomic beam. As the velocity capture range \( \Delta v = \Gamma / k \) (with \( k = 2\pi / \lambda \) and \( \lambda \) the wavelength of the laser light) of the Doppler cooling process is quite small, we use effectively curved wave fronts to increase this capture angle, as outlined in reference [15]. This can be understood as follows: While the atom is transversely slowed its longitudinal velocity remains constant, and thus it follows a curved trajectory. The Doppler shift experienced by the atom can be expressed as \( \Delta_{\text{Doppler}} = -k \cdot v \). For this quantity to remain constant as \( v \) changes direction, \( k \) needs to change direction as well. This can be achieved by matching the curvature of the atom trajectory to the curvature of the laser wave front.

To generate these curved wavefronts, two methods have been used in the literature. One is to use a very large cylindrical lens [52]. In the other method, which is the one we have chosen, the laser beam is injected at the upstream end of a set of two, almost parallel, mirrors. The injection angle is \( \beta_0 = 80 \) mrad, and the angle between the mirrors is \( \alpha \sim 1 \) mrad. The angle \( \beta \) between the laser beams and the atomic beam axis is changed by \( \alpha \) with each reflection. Thus, effectively curved wavefronts are created. The length of the mirrors is 280 mm, yielding an interaction time of 280 \( \mu s \), or 2800 \( \tau \) for 1000 m s\(^{-1}\) atoms. The width of the mirrors is 50 mm, and their thickness is 6 mm. A schematic overview of the arrangement used is shown in figure 4.4.

For flat mirrors, and in the limit of small steps \( \alpha \), the dependence of the angle \( \beta \) on the axial position \( z \) can be approximated as:

\[
\beta(z) = \sqrt{2\alpha \frac{z_{\text{max}} - z}{d}} \tag{4.2}
\]

with \( z_{\text{max}} \approx d\beta_0^2/2\alpha \) defined as \( \beta(z_{\text{max}}) = 0 \), and \( d \) is the distance between the mirrors. There are a number of considerations that come into choosing the optimal parameter set:
4.3. COLLIMATION

1. Small values for the angle $\beta$ means that the operation is relatively insensitive to the longitudinal velocity.

2. A small value for $\alpha$ means that the steps in effective detuning are small, ensuring that the atoms can follow the wavefront and be collimated.

3. As $z$ approaches $z_{\text{max}}$, the effective detuning changes more rapidly with $z$, causing atoms to go off resonance.

4. At very small angles $\beta_0$, it may be difficult to couple the laser light in.

5. For a large number of reflections $N$, the flatness of the mirrors in the $x$ or $y$ direction may become important.

To eliminate effect (3) we use mirrors with a small curvature in the $z$ direction, and making $\alpha$ a function of $z$. This is realised by slightly stressing a flat mirror using a spring, exerting pressure in the middle of the mirror, which is held at both ends. The deflection is in the order of 100 $\mu$m. If we estimate the mirror to be parabolical, we can approximate $\alpha(z) = \alpha_0 - \alpha_1 z$. To obtain a constant (transverse) deceleration, a linear dependence of $\beta$ on $z$ is required: $\beta(z) = \beta_0 - \beta_1 z$, keeping in mind that $\beta(z) > 0$ for all $z < L_{\text{mirror}}$. This gives the following simple relations between $\alpha$ and $\beta$:

$$\alpha_0 = \beta_1 \beta_0 d \quad (4.3)$$

$$\alpha_1 = \beta_1^2 d \quad (4.4)$$

Figure 4.4: Collimation arrangement
(a) One dimensional collimation

(b) Two dimensional collimation

Figure 4.5: Collimator operation

The maximum value of $\beta_1$ is determined by the maximum deceleration and the axial beam velocity. The atomic beam capture angle is determined by $\beta_0 - \beta(\text{z_{mirror}})$, where $\beta(\text{z_{mirror}}) > 0$. As mentioned (item 5), as $\beta \rightarrow 0$ the number of reflections $N$ increases rapidly, and the flatness of the mirrors becomes important. In our setup, we found that $N < 60$ gives acceptable results.

As in reference [15], two orthogonal identical sets of mirrors are used to achieve two dimensional collimation. We test the operation of the collimation section by observing the atomic beam profile, 60 cm downstream of the collimation section, using a Micro Channel Plate electron multiplier (MCP) and a phosphor screen. Light from the phosphor screen is monitored by a CCD camera and can be captured using a frame grabber. A metastable atom that hits the front face of the MCP ejects an electron that is subsequently amplified by the MCP, and accelerated onto the phosphor screen by a high voltage (5 kV). As the atom fluxes in the experiment are quite high, the MCP is used in a low gain mode, with about 500 V across the plates. In figure 4.5 the beam profile as seen by the CCD camera is displayed, with only one dimension operating (left) and with both dimensions operating (right). For the right hand image, the gain of the MCP was further reduced to avoid saturation of the camera. For a more detailed discussion of the results from the collimating section I refer to [25].
4.4. Slowing

4.4.1 Design

The design requirement for the atomic beam decelerator is that the collimated atomic beam, with a high average axial velocity \( v_{ax} = 900 \text{ ms}^{-1} \) at LN\(_2\) temperature and a slightly supersonic velocity spread, is slowed to 10 - 100 ms\(^{-1}\) in a relatively short interaction time. For the slowing stage, a liquid nitrogen cooled atom source is of paramount importance, as we shall see that the required interaction length depends linearly on the source temperature. For the deceleration of atoms we use the radiation pressure force of a laser beam, counterpropagating the atomic beam, exciting the \( 2^3S_1 \rightarrow 2^3P_2 \) transition. The maximum deceleration is \( a_{\text{max}} = \hbar k \Gamma / 2m = 4.7 \cdot 10^5 \text{ ms}^{-2} \) for full saturation. For atoms with an initial velocity of 1000 ms\(^{-1}\) the minimum stopping distance is \( v_f^2 / (2 * a_{\text{max}}) = 1.06 \text{ m} \), proportional to the initial velocity squared and hence to the source temperature. Hence, for a room temperature (300 K) source, the minimum slowing length would thus go up to \( \sim 4 \text{ m} \). However, atoms with a velocity \( v \) in the laboratory inertial frame experience a Doppler shift \(-kv\), which is \( \sim 1 \text{ GHz} \) for 1000 ms\(^{-1}\) atoms, much larger than the natural linewidth \( \Gamma / 2\pi = 1.6 \text{ MHz} \) of the transition. In order to continuously compensate for the changing Doppler shift of the decelerating atoms we have employed the Zeeman slowing technique.

With this technique, a circularly polarised laser field is used whilst the atoms are placed in a magnetic field that varies along the atomic deceleration path. Thus, the changing Doppler shift is compensated by a changing Zeeman shift in the magnetic field. The energy shift \( \Delta E \) of an atom in a magnetic sublevel \( m_j \) with respect to a magnetic field \( B \) is given by \( \Delta E = \hbar \mu_B g_j m_j B \) with \( g_j \) the Landé factor and \( \mu_B \) the Bohr magneton. The resonance condition, which gives maximum deceleration in the slowing process is given by:

\[
\Delta + \omega_e - kv = 0
\]

\[
\hbar \omega_z = \mu_B B (g_{j_e} m_e - g_g m_g)
\]

\[
g_{j_g} = 2 \quad m_g = 1
\]

\[
g_{j_e} = 1.5 \quad m_e = 2
\]

\[
\Rightarrow \hbar \omega_z = \mu_B B \quad (4.5)
\]

where \( \Delta = \omega_l - \omega_a \) is the laser detuning, \( \mu_B \) is the Bohr magneton. For a \((j = 1, m = 1) \rightarrow (j = 2, m = 2) \) transition, the Zeeman frequency shift in a magnetic field \( B \) is equal to \( \hbar \omega_z = \mu_B B / \hbar \).
CHAPTER 4. THE BRIGHT BEAM APPARATUS

If a constant deceleration \( a \) is assumed, the velocity of the atom, and hence the Doppler shift, is proportional to the square root of its position. Consequently, to maintain resonance the magnetic field \( B \) should also be proportional to the square root of the position. An arbitrary offset magnetic field \( B_0 \) can be chosen, as its effect can be compensated by the laser detuning \( \Delta = \mu_B B_0/\hbar \). Thus we arrive at the following expression for the magnetic field as a function of the position \( z \):

\[
B(z) = B_0 + B_1 \sqrt{z_{\text{magnet}} - z}
\]  

(4.6)

where \( B_1 = \hbar k v_i / \mu_B \), and \( z_{\text{magnet}} \) the interaction length. The changing magnetic field defines a decelerating, moving frame in which the deceleration remains constant. As soon as the atom comes into resonance with the slowing laser, the atomic velocity (and thus the Doppler shift) is ‘locked’ to the Zeeman shift as long as the maximum deceleration \( a_{\text{max}} \) is not exceeded. As this process is also dissipative, the longitudinal velocity distribution is bunched while the beam is being slowed. The resulting slow atom beam theoretically has a velocity spread of less than 1 ms\(^{-1}\). In practice we have found it to be less than 5 ms\(^{-1}\) while the average velocity can be tuned to anywhere between 0 and 200 ms\(^{-1}\). As the slowing process is in theory 100% effective, the full beam flux is maintained.

In our setup, we choose \( B_1 = 0.06T \) and \( B_0 = -0.03T \), and the length of the magnet \( z_{\text{magnet}} = 2.0 \) m. It is a common practice to have a Zeeman slower whose length is twice the theoretical stopping distance. There are a number of reasons for this:

- Firstly, the intensity of the slowing laser varies across the beam waist, and in some regions may not reach saturation at all.

- Secondly, due to momentum diffusion (see section 2.2) the individual atomic velocity can differ from the velocity of the decelerating inertial frame. By choosing a deceleration smaller than \( a_{\text{max}} \), the atomic velocity is ‘locked’. The atoms are then cooled, as there is a friction force around this ‘locking’ point. This friction force also counteracts heating through force fluctuations.

We use a configuration with a sign change in the magnetic field, as this yields the lowest overall magnetic fields, and also allows for additional transverse collimation of the beam in the region of the magnetic field zero [53]. An additional advantage is that the laser detuning is large, and hence the laser does not interact with the atoms after the atoms have left the magnetic field.
formers diameter | 88.8 mm
vacuum diameter   | 60 mm
former 1 length  | 1500 mm
former 2 length  | 500 mm
# layers         | 19
wire width       | 3.3 mm
wire thickness   | 1.93 mm

Table 4.1: Some parameters used for the Zeeman tuning coils in the slowing section. The number of turns as a function of position is tabulated in table 4.2. We use rectangular cross-section copper wire, wound on two formers, which consist of a cooling jacket around a stainless steel vacuum tube.

The magnetic field is generated by two tapered long solenoids. These coils are formed by winding a 2 × 3.3 mm thick rectangular cross-section wire on a double walled, water cooled vacuum tube. The inner diameter of the vacuum tube is 60 mm, and outer diameter of the cooling jacket is 88.8 mm (see table 4.1). Both magnets have 19 layers of wire, and for the first (longer) magnet the number of windings is 18 for the last layer, and 441 for the first layer; the second magnet has 151 windings in the first layer and 2 in the last one. The maximum current used is 6.5 A. The dimensions of the slowing coils are summarised in table 4.1, and the number of windings as a function of position is tabulated in table 4.2, and the calculated axial magnetic field, produced by the magnet is shown in figure 4.6.

The slowing mirror, which couples in the ~50 mm (1/e² diameter) slowing laser beam with a maximum saturation parameter of $s = I / I_{\text{sat}} = 5$, is placed after the compressor (as illustrated in figure 4.7) for the experiments which do not employ the trap. Atoms leave the compressor through the hole in the slowing mirror (diameter 3 mm). In the experimental configuration where the trap is loaded from the bright beam apparatus, the trap is placed after the slower, and the slowing mirror outside the trap. The slowing laser is a dedicated diode laser (SDL 6702-H1) with an output power of 50 mW. The slowing laser is locked to a saturated absorption setup 400-500 MHz from resonance using an AOM.
### Table 4.2: The number of turns as a function of the layer number in both Zeeman tuning coils

<table>
<thead>
<tr>
<th>layer #</th>
<th>coil 1</th>
<th>coil 2</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td># turns</td>
<td>length (mm)</td>
</tr>
<tr>
<td>1</td>
<td>441</td>
<td>1455.3</td>
</tr>
<tr>
<td>2</td>
<td>422</td>
<td>1392.6</td>
</tr>
<tr>
<td>3</td>
<td>407</td>
<td>1343.1</td>
</tr>
<tr>
<td>4</td>
<td>390</td>
<td>1287</td>
</tr>
<tr>
<td>5</td>
<td>370</td>
<td>1221</td>
</tr>
<tr>
<td>6</td>
<td>352</td>
<td>1161.6</td>
</tr>
<tr>
<td>7</td>
<td>329</td>
<td>1085.7</td>
</tr>
<tr>
<td>8</td>
<td>308</td>
<td>1016.4</td>
</tr>
<tr>
<td>9</td>
<td>286</td>
<td>943.8</td>
</tr>
<tr>
<td>10</td>
<td>262</td>
<td>864.6</td>
</tr>
<tr>
<td>11</td>
<td>239</td>
<td>788.7</td>
</tr>
<tr>
<td>12</td>
<td>215</td>
<td>709.5</td>
</tr>
<tr>
<td>13</td>
<td>188</td>
<td>620.4</td>
</tr>
<tr>
<td>14</td>
<td>161</td>
<td>531.3</td>
</tr>
<tr>
<td>15</td>
<td>135</td>
<td>445.5</td>
</tr>
<tr>
<td>16</td>
<td>109</td>
<td>359.7</td>
</tr>
<tr>
<td>17</td>
<td>78</td>
<td>257.4</td>
</tr>
<tr>
<td>18</td>
<td>47</td>
<td>155.1</td>
</tr>
<tr>
<td>19</td>
<td>18</td>
<td>59.4</td>
</tr>
<tr>
<td>total</td>
<td>4757</td>
<td>1500</td>
</tr>
</tbody>
</table>
4.4. SLOWING

4.4.2 Deceleration results

The performance of the deceleration stage was tested using the experimental setup shown in figure 4.7. A mechanical chopper, placed 360 mm downstream of the source and just downstream of the collimator was used for the time of flight measurement of the atomic velocity. The time-of-flight technique has the advantage that it does not require an extra probe laser and consequently does not rely on observation of fluorescence, which for the 1083 nm wavelength is very difficult, even with a dedicated camera. During these measurements the slowing mirror, used to couple the slowing laser into the slower is placed just after the

Figure 4.6: The calculated axial magnetic field in the Zeeman slower

Figure 4.7: Differential time-of-flight setup
Chapter 4. The Bright Beam Apparatus

UV photons
detector 2
Unslowed
(detector 1
Slow atoms
At=4 ms

counts

8000
6000
4000
2000
0

0 100 200 300 400
channel

Figure 4.8: Time-of-flight spectra as recorded by the two detectors. Each TOF-channel is 50 μs. The axial distance between the detectors is 400 mm. Consequently a difference in time-of-flight of 4 ms indicates a beam velocity of 100 ms\(^{-1}\).

slower. We use two channeltrons as detectors in this experiment, as depicted in figure 4.7. The first channeltron is mounted adjacent to the slowing mirror and has a positive voltage (200 V) at its front and 3000 V at the back. The whole channeltron box is at the potential of 100 V with respect to the chamber. The higher positive voltage at the front of the channeltron accelerates electrons towards the detector. This channeltron detects the electrons generated by He\(^*\) atoms that are not well enough collimated to go through the 3 mm hole in the center of the slowing mirror, and which strike the mirror. The flight time at detector 1 determines the average velocity for the atoms in the slower. Those atoms that pass through the hole in the slowing mirror, are detected by the second channeltron placed after the slowing mirror. The front end of this channeltron is at zero potential, and the back is at +3000 V. The second detector is placed 400 mm downstream from the slowing mirror. Figure 4.8 shows the time-of-flight spectra recorded using the two detectors. The difference in time-of-flight yields a measurement of the velocity of the atoms after the slower. The atomic velocity \(v_z = \Delta z / \Delta t\) deduced in this manner is less than 100 ms\(^{-1}\).
4.5 Compression

4.5.1 Introduction

After the first collimating section, we end up with a parallel, high flux atomic beam of rectangular profile with a relatively uniform peak distribution and a width of $\sim 20$ mm. The slower that follows the first collimating section narrows the longitudinal velocity distribution, but inevitably increases the transverse velocity spread through diffusive heating. The atomic beam at the end of the slower has a diameter of $\sim 50$ mm.

The requirement for the bright beam is a high phase space density of the beam. To decrease the beam width, we have to create a two dimensional position sensitive force or a two dimensional trap for neutral atoms. Several options have been considered:

- The atoms can be forced to move towards the axis by the position dependent AC-Stark shift generated by a high power red detuned laser beam counterpropagating the atomic beam. A disadvantage of this method is the large amount of laser power required. Also, the phase space density is not increased using this method.

- Another possibility is to use a hexapole magnetic lens, which focusses an atomic beam of polarised atoms, as demonstrated for a laser slowed caesium beam in [54]. In the case of a non polarized atomic beam, as ours is, this would only focus one third of all atoms.

- We have chosen the method which is suitable for a nonpolarised atomic beam, requiring only a modest amount of laser power. We are using a position dependent light pressure force in an inhomogeneous magnetic field, taking advantage of the Zeeman effect, as first outlined in reference [55] and used in references [13, 14, 15]. It uses a two dimensional magnetic quadrupole field, combined with two sets of counter propagating, circularly polarised laser beams. This essentially constitutes a two-dimensional magneto-optic trap (MOT). The two-dimensional quadrupole magnetic field is shown schematically in figure 4.9, as well as the definitions of the spatial axes.

The basic operation of a one dimensional MOT for an atom with a $J = 0 \rightarrow J = 1$ transition is illustrated in figure 4.10. The scheme can easily be generalised to two
CHAPTER 4. THE BRIGHT BEAM APPARATUS

Figure 4.9: The two-dimensional quadrupole field and the axis definitions

Figure 4.10: The Zeeman tuned transition frequencies as a function of $x$ in a simplified picture. The $\sigma^+$ polarized laser beam from the right has a small red detuning. Atoms at the position $x_1$ are more resonant with the $\sigma_+$ laser beam, coming from the left and will therefore, be forced towards the beam axis by radiation pressure. The $\sigma^-$ polarized laser beam from the right is most resonant with atoms at position $x_2$. All atoms are thus pushed towards $x = 0$.

or three dimensions and larger angular momenta of the atom. For a small laser detuning $\Delta \sim -\Gamma$, the position sensitive, 'restoring' force created is proportional to the distance from the axis of the magnetic quadrupole.

When the atom with ground state $g$ and excited state $e$ is illuminated with either $\sigma^+$, $\pi$, or $\sigma^-$ polarized light with respect to the magnetic field, transitions
with $\Delta m = +1, 0$ or $-1$ are induced. The transition frequency is then shifted by

$$\Delta_{\text{mag}} = \frac{\mu_B B [g_{j_e}(m_g + \Delta m) - g_{j_g}m_g]}{\hbar}$$

(4.7)

By introducing a magnetic field in the direction of the $k$-vector of the light we can compensate for the laser detuning in circularly polarized laser light. Therefore, the radiation pressure force depends on the position of the atom in the inhomogeneous magnetic field. In a quadrupole magnetic field, the magnetic field is zero on the axis and increases linearly with the distance from the axis with a gradient $G$.

We now illuminate the atom with a laser field with $\sigma^+$ light running in the $+x$ direction and $\sigma^-$ light running in the $-x$ direction. In a simplified picture we assume that in a weak magnetic field the frequency difference between the transitions with identical $\Delta m$ to be small, which is true when $|g_{j_e} - g_{j_g}| \ll g_{j_e}$ (see equation 4.7). In figure 4.10 the simplified energy diagram is shown as a function of $x$. The laser frequency is detuned to the red of the resonance line by an amount $\Delta$. At position $x_1$, the laser detuning is exactly compensated for by Zeeman shift for the $\Delta m = +1$ transition. Consequently an atom at position $x_1$ is more resonant with the $\sigma^+$ light, which is running in the $+x$ direction and will thus experience a force towards the axis. The reverse is true for an atom at position $x_2 = -x_1$. For a long interaction time all atoms will carry out a damped oscillatory motion around and towards the quadrupole axis [55]. The oscillation can be damped rapidly by ramping up the magnetic field gradient $G$ as the atoms traverse the interaction region.

The magnetic field gradient at the beginning of the compressor section $G(z = 0)$ is determined by the condition that the Zeeman shift for the outermost atoms should equal the laser detuning, modified by the Doppler shift due to the transverse velocity. In the current setup, the laser detuning in the compressor stage is relatively large (typically $\Delta \sim -10\Gamma$), to compensate for the large transverse velocity spread of atoms coming out of the Zeeman slower. This transverse velocity spread can be estimated from the number of photon recoils ($n_p \sim 10^4$) required to slow the atoms in the Zeeman slower, and the assumption that the heating is proportional to the square root of this number. So estimated, the transverse velocity spread is $\sqrt{(n_p)\hbar k/m} \approx 9\text{ms}^{-1}$. The initial transverse velocity spread ($\sim 5\text{ms}^{-1}$) should be added in quadrature. Ideally, the magnetic field gradient $G(z)$ should increase linearly from a small value to 40 G/cm. In the first half of
the compression stage this increasing gradient compensates for the Doppler shift introduced by the atoms starting to move away from the laser beam that pushes them towards the quadrupole axis. In the second half of the compression stage, the large gradient ensures tight confinement of the atoms on the axis.

### 4.5.2 Simulation

The trajectories of the atoms in the various stages of the beam brightener are simulated in two dimensions using a simple integration routine. This program has been used in the past to simulate a similar setup for neon [49] and has been adapted to helium and expanded to include the calculated magnetic field (see section 4.5.4). The operation of the program can be outlined as follows: Each atom gets an initial velocity, position and magnetic substate. The scattering rate from each of the laser beams is calculated, taking detuning, Doppler shift, Zeeman shift and (cross) saturation into account. From this we obtain the resulting force. The motion of the atom is adjusted, and a momentum kick, proportional to the total scattering rate but in a random direction is added to simulate momentum diffusion. Optical pumping is also taken into account. The motion of the atom is then propagated for the next integration step.

A three dimensional simulation, accounting for atoms that are not on either the $x$ or the $y$ axis has also been developed for neon [49] and adapted to helium. In the compressor calculations, this presents a significant complication. The magnetic field, which defines the quantisation axis for the magnetic sublevels, is not in the direction of the laser beams. The laser beams are circularly polarized ($\sigma^+ / \sigma^-$) with respect to their wave vectors. Using rotation operators we can express the circular laser polarizations as linear superpositions of $\pi$, $\sigma^+$ and $\sigma^-$ light with respect to the local magnetic field direction. The absorption rate of each of these polarisation components depends on the Zeeman resonance condition. In the simulation, the detuning in equation 2.3 is modified with the Zeeman detuning $\Delta_{\text{mag}}$. All laser beams are expressed in their polarization components with respect to the magnetic field, each with a different Zeeman detuning. The effective force for each polarization component of each laser beam is again calculated using equation 2.3. The three dimensional simulation is very slow in comparison, and yields little insight that cannot be obtained from the two-dimensional simulation, adjusting for cross-saturation by always choosing the laser power well below saturation.
4.5. COMPRESSION

In the simulations, as well as in tests using the compressor as a lens for fast atoms, the compressor generally appears to operate with reasonable axial symmetry.

4.5.3 Design

To obtain rapid damping of the "oscillatory" motion in the trap, and thus a short interaction time, it is advantageous to use an increasing magnetic field gradient $G(z) = dB_{(x,y)}/dz$ as the atoms traverse the laser field in the $z$ direction. This can be realised by using the fringe fields of the magnets: the quadrupole strength decreases as $z - z_{magnet}$ increases.

The long lifetime of the $^3P$ state of helium (100 ns) dictates a long interaction length to obtain fully damped transverse motion. Simulations show that a 70 ms$^{-1}$, 50 mm diameter atomic beam can be compressed in 100 mm. This requires an area of 50 cm$^2$ to be illuminated homogeneously with the correct laser polarisation from four sides. We use an arrangement of polarisation preserving [56], in-vacuum mirrors as proposed by Driessen [57]. A schematic overview is shown in figure 4.11. A single circularly polarised input laser beam is reflected 32 times to create the desired laser fields, thus greatly reducing the laser power required for reliable operation. The assembly uses a rooftop arrangement to

---

Figure 4.11: The mirror arrangement used in the compression stage Views: a) 3D view, b) view along the atomic beam axis
flip from $\sigma^+$ to $\sigma^-$ polarisation for the retroreflected beam. To compensate for any phase error $\Delta \Phi$ between the $s$- and $p$-components of the laser field, the reflections alternate between "horizontal" and "vertical". This way, the maximum phase error accumulated during the 32 reflections is $2\Delta \Phi$. We have verified that the light exiting the mirror arrangement after 32 reflections is indeed circularly polarised to the measurement accuracy. An added advantage of this mirror arrangement is that the magnets can be placed around part of the interaction region, and in this way they do not block any incoming laser beam.

The relative alignment of all mirrors in the assembly is very critical, as the laser beam travels more than a metre inside the assembly, and there are very many degrees of freedom. We have taken an incremental approach to this alignment problem: First we ensure that the various mirror pairs are at right angles. This is verified by illuminating the apex with a large area, well collimated laser beam and observing the continuity of the fringes in a Shearing interferometer (Melles Griot 09-SPM-003) placed in the reflected beam, across the reflection on the apex. If the angle in the apex is not exactly a right angle, the two halves of the reflected beam will have different $k$-vectors, leading to a small fringe spacing in the interferometer. In figure 4.12 we show such an image in the interferometer for proper alignment (centre) and for a small misalignment (left and right). Using this method, the right angles are accurate to better than 1 mrad. Then, the relative alignment of the mirror pairs is adjusted by observing the output beam and the degree to which the volume between the mirrors is filled with laser light. This relative alignment was found to be somewhat less critical. All alignment can be done outside the vacuum. Then, the entire assembly is placed in the vacuum chamber, and the input laser beam is aligned accordingly.

4.5.4 The quadrupole magnetic field

The quadrupole magnetic field is formed by four ($50 \times 25 \times 100$ mm) NiFeB magnets. They are placed 110 mm away from the atomic beam axis and $\sim 50$ mm downstream from the center of the interaction region. The field gradient $dB_{x,y}/dx\{x,y\}$ thus obtained ranges from $1 \cdot 10^{-2}$ T/m at the beginning of the interaction region to $40 \cdot 10^{-2}$ T/m at the end of the interaction region. To reduce the initial field gradient, an additional set of four ($50 \times 12 \times 100$ mm) magnets can be placed $\sim 100$ mm upstream of the interaction region, and up to 200 mm away from the atomic beam axis.
Figure 4.12: The method for inspecting the correct alignment of a rooftop mirror pair. The vertical lines are the reflections from the front and back face of the Shearing plate of the apex between the mirrors. In the area between these lines, parts of laser beam following different paths interfere, giving a sensitive indication of the angle difference between these parts, and hence of the angle in the apex. The horizontal line is a visual aid on the interferometer, indicating the direction of the fringes for a collimated laser beam.

The magnetic field can be calculated from first principles, assuming that all magnets are sufficiently far apart so that they do not influence each other’s fields, but can just be added linearly. For each point in space, we calculate the distance from each magnet, and its magnetic field at that point. The magnet can have any orientation. The fields from all the magnets are then added. To obtain the magnetic field from one magnet, we use the magnetic potential $V_m$, defined as $B = -\nabla V_m$. Of course, this is only a valid concept in the absence of currents or magnetically polarisable materials. The magnetic potential can be found from a simple 3D integration of the magnetic field strength. For a ring or rectangular magnet, one dimension of the integration can be carried out analytically (using Maple, a computer algebra package). Inserting this in the integration significantly speeds up the integration, which is done numerically using a NAG routine (D01AHF), a commercially available numerical routine developed by the Numerical Algorithms Group (NAG). The results from the magnetic field calculations can easily be inserted in the atomic trajectory simulation.

The overall shape of the magnetic field, as measured using a Hall probe, is in good agreement with the calculations. However the magnitude is significantly larger than expected from the magnet specifications, which is due to the larger magnetisation of the magnets used.
4.5.5 Operation

To obtain a low divergence of the output beam from a two-dimensional MOT, one might expect large detunings to be optimal, as has been found in three dimensional MOTs and other atomic species [32]. In contrast to these findings, for our atomic beam parameters a small laser detuning is desirable ($\Delta \sim -\Gamma/2$). This can be understood as follows. For a two-dimensional MOT, it is more important to exert a large force on the atoms during a limited interaction time then it is to reduce the heating rate when the atoms are at zero velocity. Furthermore, in the case of helium, polarisation gradient cooling, which is responsible for the achievement of lower (transverse) temperatures at higher detunings, does achieve a significantly lower temperature, as the recoil limit is very close to the Doppler limit. Hence, to obtain a low divergence of the output atomic beam from the two dimensional MOT, a small laser detuning is more effective.

However, to ensure a large capture range at the input side of the laser field, this translates to a very small magnetic field gradient. These conditions lead to two important complications in the operation of the compression stage:

1. The increase of the gradient $dG(z)/dz$ is substantial. Maxwell's equations ($\nabla \cdot \vec{B} = 0$) demand that this is paired with a substantial axial magnetic field $B_z$. This axial magnetic field effectively obstructs the optical pumping effect needed for an efficient restoring force. In practice, this means that the available restoring force is much smaller than could be deduced from the simulations, which do not include axial magnetic fields. Another consequence is that the gradient $G(z)$ needs to have a weaker dependence on $z$. Consequently, the required interaction time becomes larger.

2. The small laser detuning also causes problems with the capture of the strongly divergent atomic beam that exits the slowing stage. A larger laser detuning $\Delta \approx 20$ MHz in the compression stage can rectify this problem.

These complications suggest the use of two separate stages in the compression:

1. A focusing stage, with a large laser detuning, and consequently a relatively large initial magnetic field gradient, enabling the capture of the full transverse velocity distribution exiting the slower.

2. A recollimation stage, with a smaller laser detuning to ensure sufficient collimation of the final atomic beam.
In the present work, we have used this two stage scheme to obtain a high brightness atomic beam. The atomic beam is focused on the hole in the mirror that is used to reflect the slowing laser beam on axis. After this mirror, a two-dimensional recollimation stage, consisting of a 20 mm long polarisation gradient optical molasses is used. The residual magnetic field in this region was compensated by using a number of extra permanent magnets outside the vacuum system. Although no compression is achieved in this stage, this serves to create a $\sim 5$ mm diameter, well collimated atomic beam with a beam flux of in excess of $10^{10} \text{s}^{-1}$ at an axial velocity of $\sim 100 \text{ms}^{-1}$. In the current magnetic field configuration, this is the only practical beam velocity as other beam velocities are not focussed on the hole in the mirror and hence suffer a substantial loss in beam flux. Figure 4.13 shows the atomic beam profile 300 mm downstream of the recollimation section, as measured using the MCP/phosphor screen/CCD camera combination, as detailed in section 4.3. To measure the beam flux, a fine stainless steel mesh, with an open fraction of approximately one half, is mounted a few millimetres upstream of the MCP. Whilst the image on the phosphor is not degraded, half of the atoms hit the mesh, and eject an electron. The current from this mesh is in the nA regime, and can easily be measured using a standard electrometer (Keithley 610B). The other half of the atoms is transmitted by the mesh and is detected by the MCP. The high voltages used in the MCP ensure that no space charge is built up around the mesh, as the electrons ejected from the mesh are immediately accelerated towards the high voltage electrodes.

4.5.6 Conclusions and suggested improvements

Although the compression stage, with additional collimation, currently yields a well collimated and intense beam, the overall yield of the compression scheme is not optimal. The flux of the collimated atomic beam after the collimation stage is $\sim 10^{12} \text{s}^{-1}$. This in contrast with the compressed beam flux of $\sim 10^{10} \text{s}^{-1}$. Consequently, only 1\% of the atoms leaving the collimation stage is captured in the compressed, slow beam. The flux after the slower and the compressor is $\sim 2 \cdot 10^{11} \text{cm}^{-2}\text{s}^{-1}$, and the corresponding density is $\sim 2 \cdot 10^7 \text{cm}^{-2}$. Given a 100\% capture efficiency and no other losses, these numbers could be a factor of 100 larger.

Still, the density is equal to that a few centimetres away from the discharge source, and is now in a collimated, monochromatic, slow beam that propagates
CHAPTER 4. THE BRIGHT BEAM APPARATUS

Figure 4.13: Beam profile of the slow atomic beam as measured by the MCP/phosphor screen/CCD camera combination (left) and a 3D representation of this beam profile (right)

The grey scale value of the image is linearly related to the beam flux.

in a background-free environment. It is anticipated that the current setup can be used for a range of experiments that do not require a very high flux.

However, the potential for higher atomic beam fluxes suggests further research to increase this flux. A number of possibilities exist to study and improve the operation of the compression stage:

- Simulations suggest that the operation can be greatly improved by using a "pre-focussing" stage, in which a larger laser detuning enables the capture of a large range of transverse velocities. Subsequently, the detuning in the main compression stage can be reduced, and hence a lower divergence atomic beam can be obtained. In figure 4.14 the atomic trajectories in such a two-stage compressor are displayed.

- So far, the alignment in the compression mirror array has proven to be a stumbling block. A critical parameter is the relative alignment of the "sets" of perpendicular mirrors. However, there are no convenient checks for these angles while the mirror sets are not bolted together. A further reference plane, enabling us to rest the mirror array on its side may prove to give
Figure 4.14: Atomic trajectories in velocity space and real space for a two-stage compression setup. The first 5 cm act as a focussing lens. In a 10 cm drift region the atomic beam converges, and the beam is compressed in the following 10 cm. In this simulation, the laser intensity $I = I_0$, and the laser detuning $\Delta = -5\Gamma$. This large detuning results in a large velocity capture range, but yields imperfect collimation at the exit of the compressor. This may be remedied by adding a further re-collimation stage.

A valuable insight in this problem. Currently, such a reference plane is being built into the mirror holder.

- A limiting factor for the obtainable density is Penning ionisation arising from collisions between a $^3S$ atom and a laser excited $^3P$ atom, as detailed in chapter 6. Consequently, if the compression stage works too well, the beam will autodestruct. The importance of this mechanism could be studied using a channeltron detector, set up to detect the resulting ions close to the compression stage.

- A major obstacle for optimisation of the compression stage is the fact that the detected beam has to pass through the hole in the slowing mirror, as the MCP detector is not transparent for the slowing laser light. Optical detection methods (absorption, emission at 588 nm) might circumvent this problem, and give a clear idea of where all the atoms end up.
Chapter 5

Generation of 389 nm light

5.1 Motivation

This chapter reports on the experimental realisation of an apparatus designed to generate light at 389 nm to excite the $2^3S_1 \rightarrow 3^3P$ transition in metastable He (figure 1.1.). Most of the work described in this thesis is done using the $2^3S \rightarrow 2^3P_2$ transition, from the metastable state, and this is used to manipulate the atomic trajectories in the bright atomic beam apparatus. We have planned several further experiments using the bright beam source that utilize the 389 nm transition in He*.

This transition is particularly interesting due to the fact that quantum effects are more pronounced. The recoil velocity for this transition $v_{\text{rec}} = \frac{\hbar k}{M} = 25.6 \text{ cms}^{-1}$ is almost three times larger than the (already large) recoil velocity for the 1083 nm transition, which is $v_{\text{rec}} = 9.20 \text{ cms}^{-1}$. In the planned VSCPT experiment (see section 2.4) the splitting angle is proportional to the recoil velocity, and this would be three times larger for the 389 nm transition. Not only does this have obvious advantages for the detection, but also the enclosed area of an atom interferometer based on this scheme is proportionally larger, increasing the sensitivity.

The ultimate application of 389 nm light in this thesis has been to modify cold atom collisions in the trap experiments, described in chapter 6. Enhanced ionisation in the presence of red detuned light is an interesting example of the dynamics that may occur in ultracold samples of atoms. Conversely for blue detuned light, it is expected that the ionisation rate will be suppressed, as has been observed in xenon [58, 59]. The 389 nm laser may provide an excellent tool for studying cold atom collision dynamics, and to increase the trap density. We
use a Titanium Sapphire laser, frequency doubled in an LBO crystal cavity, to generate the 389 nm light used to perform these experiments.

5.2 Frequency doubling

Optical frequency conversion was first demonstrated in 1961 by Franken, Hill, Peters and Weinreicht [60]. The conversion efficiency achieved in this, first experiment was $10^{-8}$. Since then, theoretical and experimental progress in conjunction with technological advances, have raised the efficiency of the doubling process to more than 0.8 for some systems.

When an electromagnetic wave passes through a dielectric material, the valence electrons are displaced from their normal orbits. As a result, temporary dipoles are formed in the material, creating a polarisation wave which radiates an electromagnetic field. For low electromagnetic field strengths, the polarisation wave mimics the incoming wave. This linear dependence of the material polarisation on the electric field, can be used to explain phenomena such as reflection, absorption, normal and anomalous dispersion, and forms the field of linear optics. At higher field strengths the dipole response is distorted, and becomes dependent on the higher order powers of the time dependent field. The distorted re-radiated wave contains harmonics of the fundamental frequency. The induced polarisation, $P$, can be expanded in a convergent power series [61]:

$$P = \chi^{(1)}E + \chi^{(2)}E^2 + \chi^{(2)}EB + \chi^{(3)}E^3 \ldots \tag{5.1}$$

where $\chi^{(1)}$ is the linear susceptibility, and $\chi^{(2)}$ and $\chi^{(3)}$ are (generally weaker) higher order nonlinearities in the dielectric response. In this chapter we consider optical $\chi^{(2)}$ processes. These are three photon processes where two low frequency photons are converted to one high energy photon (upconversion), or the other way around (downconversion). Second harmonic generation is an example of upconversion. The low frequency field is known as the fundamental, and the high frequency field is the second harmonic.

Both the energy and momentum of the interacting photons must be conserved, which can be expressed respectively as:

$$\hbar \omega + \hbar \omega = 2\hbar \omega \tag{5.2}$$

$$\vec{k}_\omega + \vec{k}_\omega = \vec{k}_{2\omega} \tag{5.3}$$
The latter is the condition of phase matching. We can define the phase mismatch as:

$$\Delta k = 2k_\omega - k_{2\omega} = 2\frac{n_1\omega}{c} - 2\frac{n_2\omega}{c}$$

(5.4)

In $\chi^{(2)}$ materials, the relationship between the wave vector $k$ and refractive index of the material, $n$ is [61]:

$$k(\nu, \Pi) = \frac{\omega}{c} n(\omega, \Pi),$$

(5.5)

Both wave vector and index of refraction are functions of the frequency $\omega$ and the polarisation $\Pi$. The polarisation wave, generated in the medium by the incident wave, radiates light with different frequencies. These experience different indexes of refraction, and consequently, different light speeds. They are generated at different places with the phase distribution of the local incident wave, but with different speed of propagation. The generated wave at one position can therefore interfere with that from another position. Thus the conversion efficiency of the fundamental light to the higher harmonics depends on the interaction length and the difference in the refractive indices. The variation of the field amplitude of the second harmonic as a function of the field amplitude of the fundamental light can be expressed as:

$$E_{2\omega} \approx E_\omega^2 \sin \frac{1}{2} \Delta k z$$

(5.6)

where $z$ is the interaction length. The fact that the refractive index in crystalline materials is polarisation dependent allows the possibility to achieve phase matching. The term phase matching means that the polarisation wave and the higher frequency, generated wave are in phase over the whole interaction region. Second harmonic waves, generated at different points in the material, will then interfere constructively. There are currently two methods for achieving phase matching in a crystalline medium:

1. Type I phase matching. The refractive indices can be made the same when the two fundamental fields are of the same polarisation, and the second harmonic field is orthogonal. This can be achieved by orienting the crystal to a certain angle, setting the crystal to a certain temperature, or both. This is known as critical phase matching. If the refractive indices are matched for light propagating at $90^\circ$ to the optical axis the crystal is said to be
5.2. FREQUENCY DOUBLING

noncritically phase matched. It is noncritical in the sense that light can propagate in any direction in the \( zy \) plane, where \( z \) is the plane of the optical axis.

2. Type II phase matching can, like type I phase matching, be achieved by orienting the crystal to a certain angle, setting the crystal to a certain temperature, or both. In this way the refractive indices can be made the same when two fundamental fields are orthogonal to each other, and the second harmonic field is in the same plane as one of the fundamental fields.

We refer to the direction in which the refractive index is independent of the direction of polarisation as the optical axis. Waves plane polarised perpendicular to the plane formed by the optical axis and the direction of propagation, are referred to as “ordinary” waves (subscript ‘o’). Waves that are plane polarised parallel to this plane are referred to as “extraordinary” waves (subscript ‘e’). The following two equations describe Type I and type II phase matching.

Type I phase matching:

\[ n_e(2\omega) = n_o(\omega), \quad (5.7) \]

Type II phase matching:

\[ n_e(2\omega) = \frac{1}{2}(n_o(\omega) + n_e(\omega)) \quad (5.8) \]

G.D. Boyd and D.A. Kleinman have developed the theory of second harmonic generation, using Gaussian beams, in 1968 [62]. We use the results of their theory to design our setup. The second harmonic coefficient \( \gamma_{\text{shg}} \) is given by [62]:

\[ \gamma_{\text{shg}} \equiv \frac{P_2}{P_1^2} = \frac{2\omega_1^2 d_{\text{eff}} k_1 l}{\pi n^3 \varepsilon_0 c^3} h(B, \psi) \quad (5.9) \]

where:
CHAPTER 5. GENERATION OF 389 NM LIGHT

\( P_{1(2)} \) is the power at the fundamental (second harmonic) frequency

\( n \) is the refractive index of the crystal,

\( \varepsilon_0 \) the dielectric constant of the vacuum,

\( c \) the vacuum speed of light,

\( l \) length of crystal,

\( \rho \) walk off angle (angle between the two Poynting vectors of the fundamental and second harmonic), defined by \( \tan(\theta_m + \rho) = (n_o/n_e)^2 \tan \theta_m \)

\( \theta_m \) the phase matching angle,

\( d_{eff} \) effective nonlinear constant of the crystal,

\( k_{1(2)} \) wave vector of the fundamental (second harmonic) in the crystal,

\( w_0 \) minimum beam waist of the fundamental beam,

\( b = k_1 w_0^2 \) confocal parameter,

\( \psi = l/b \) focusing parameter

\( h \) is the focusing factor which, for \( 4B^2 > \psi > 6/B^2 \) is given by:

\[
h(B, \psi) = \frac{l_a}{l} \arctan \psi
\]  

(5.10)

where \( l_a \) is the aperture length:

\[
l_a = \frac{\sqrt{\pi}w_0}{\rho}
\]  

(5.11)

and \( B \) is given as:

\[
B = \frac{1}{2} \rho \sqrt{lk_1}
\]  

(5.12)

This leads to the following expression for \( \gamma_{shg}(l, w_0) \):

\[
\gamma_{shg}(l, w_0) = \frac{2k_1^2 d^2 w_0}{\sqrt{\pi} n^5 \varepsilon_0 c \rho} \tan^{-1}\left(\frac{l}{k_1 w_0^2}\right)
\]  

(5.13)

In figure 5.1 \( \gamma_{shg} \) is plotted versus \( l \) and \( w_0 \) for the parameters in our experiment \( (2\pi c/\omega_l = 778 \text{ nm}, \text{ LBO crystal}) \). Figure 5.2 shows a number of cuts through figure 5.1 for various crystal lengths. The figure shows that efficient second harmonic generation strongly depends on the size of the focus, for any given length \( l \) of the crystal.

5.3 Setup

The generation of a coherent, continuous laser field in the ultraviolet has been demonstrated using frequency doubling, e.g., references [62, 63, 64]. We have
Figure 5.1: $\gamma_{\text{sh}}$ as a function of the crystal length $l$ and laser beam waist $w_0$ for the parameters in our experiment ($\lambda = 778$ nm, LBO crystal).

more or less followed the arrangement of the Stony Brook group [65] with changes we considered appropriate for our experiment. The scheme is straightforward: it consists of a build-up cavity for the fundamental light in which the doubling, LBO, crystal is placed at the beam waist of the cavity. The cavity is locked to the laser frequency. As we shall see, we have obtained an output of 30 mW of blue light at 389 nm, which was enough for the purposes of our experiment. There is considerable room for enhancement of the output power, and we have analysed some ways it may be achieved in detail, if needed in the future.

The second harmonic generation (SHG) doubling cavity used in this experiment consists of two spherical and two flat mirrors. Having a flat input coupler allowed easier alignment, as the position of the beam on the flat mirror is not as critical, and this makes the adjustment of the beam waist much easier than in a cavity with spherical mirrors only.

Section 5.4 describes the choice of crystal, section 5.5 outlines the decision to use a traveling wave cavity, and section 5.6 describes the cavity locking technique.
Figure 5.2: A number of cuts through figure 5.1, which shows a second harmonic generation coefficient $\gamma_{\text{shg}}$ for various crystal lengths (5 mm, 10 mm, 15 mm, 18 mm, 20 mm, 25 mm and 30 mm respectively) for the parameters in our experiment ($\lambda = 778$ nm, LBO crystal)

5.4 The Crystal

There are three nonlinear crystals which are suitable for frequency doubling of a single frequency Ti:sapphire laser operating at 778 nm: Lithium Triborate (LiB$_3$O$_5$ or LBO), Beta-Barium Borate ($\beta$BaB$_2$O$_4$ or BBO) and Lithium Iodate (LiIO$_3$). The properties of these three crystals are compared in table 5.1. Second harmonic coefficients $\gamma_{\text{sh}}$ were calculated using equation 5.2 assuming a crystal length of 18 mm. The phase matching angle and walk-off angle were calculated using the Sellmeier equations for LBO [66], for BBO [67] and for LiIO$_3$ [68]. All of them have to be critically phase matched because they have relatively small nonlinearities. Some of the relative advantages of these crystals are:

LiIO$_3$

- LiIO$_3$ has the highest nonlinearity (see table 5.1) which is a reason to use it for frequency doubling of a Ti:Sapph laser [63]. However recently it has been shown that the nonlinear coefficient for second harmonic generation
using LiIO$_3$ is only 58% of the value previously determined by parametric fluorescence [69].

- LiIO$_3$ has a significantly lower damage threshold than the other two crystals and should not be used in the case of tight focusing and high intensities
- LiIO$_3$ is hygroscopic and requires a dry nitrogen purge during use.

**LBO and BBO**

- The second harmonic conversion coefficient of BBO is lower than LiIO$_3$ and does not offer any significant advantages.
- The second harmonic conversion coefficient of LBO and LiIO$_3$ are comparable.
- LBO has a much smaller walk-off angle than the other two crystals, and this turns out to be more important in the case of strong focusing. The smaller walk-off angle means a smaller astigmatism of the second harmonic beam.
- The surface quality of nonhygroscopic materials like LBO can be maintained to a higher quality. This is important as the surface quality determines the loss in the build up cavity and therefore the doubling efficiency.
- Additionaly LBO has a very high damage threshold (the highest among all the commonly used inorganic NLO crystals) [70]. That makes it one of the best candidates for high average power SHG.

All these reasons should ensure the better performance of LBO, in comparison to BBO and LiIO$_3$, and were therefore deciding factors for the choice of LBO as the nonlinear element in our experiment. LBO is a negative biaxial, orthorombic crystal, built up of a continuous network of B$_3$O$_7$ groups with lithium cations locating in the interstices. The compact network of B$_3$O$_7$ groups makes the LBO crystal almost free of inclusions [70].

In order to maximize the transmission of the fundamental wave through the crystal, the crystal is usually cut at Brewster’s angle, but as the generated second harmonic light is polarised perpendicular to the incoming beam (type I phase matching) the blue light leaving the crystal will be partially reflected. As given
CHAPTER 5. GENERATION OF 389 NM LIGHT

<table>
<thead>
<tr>
<th>crystal</th>
<th>LBO</th>
<th>BBO</th>
<th>LiIO₃</th>
<th>reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>(n_0)</td>
<td>1.61</td>
<td>1.66</td>
<td>1.86</td>
<td>[70]</td>
</tr>
<tr>
<td>phase matching angle</td>
<td>33.73°</td>
<td>30.02°</td>
<td>43.2°</td>
<td>[70]</td>
</tr>
<tr>
<td>(d_{\text{eff}}) (pm/V)</td>
<td>1</td>
<td>1.9</td>
<td>3.04</td>
<td>[69, 63, 70]</td>
</tr>
<tr>
<td>walk off angle (mrad)</td>
<td>17</td>
<td>69</td>
<td>87</td>
<td>[70]</td>
</tr>
<tr>
<td>damage threshold (GW cm⁻²)</td>
<td>2.5</td>
<td>1.5</td>
<td>0.5</td>
<td>[71]</td>
</tr>
<tr>
<td>(\gamma_{\text{sh}}(10^{-4}W^{-1}))</td>
<td>1.52</td>
<td>1.33</td>
<td>1.91</td>
<td>equation 5.2</td>
</tr>
</tbody>
</table>

Table 5.1: Properties of LBO, BBO and LiIO₃ at 778 nm

in (eq. 21.a [72]) the reflectivity of the light polarised perpendicular to the plane of incidence is given by

\[
R_\perp = \frac{\sin^2(\Theta_i - \Theta_t)}{\sin^2(\Theta_i + \Theta_t)}
\]

so for \(n = 1.61\) and \(\tan \Theta_i = 1/n\), the reflectivity of the second harmonic light intensity on the interface crystal to air is

\[
R_\perp = 0.196.
\]

This means that in the case of a low or zero astigmatism cavity, a Brewster cut crystal reduces the doubling efficiency by nearly 20% (the fundamental is resonant, the blue light passes the crystal only once). To avoid any further reduction, according to references [73, 69] the doubling cavity has to be designed to give a circular waist inside the crystal. Those were the reasons for having the crystal with surfaces perpendicular to the incoming beam and anti-reflective (AR) coated, with a coating with a high enough damage threshold and optical quality. [70].

5.5 The Cavity

In deciding on the cavity design, a choice needs to be made between singly and doubly resonant cavities, and standing and traveling wave cavities. The cavity has to be designed to give the optimal waist size in the crystal, and the incoming Ti:Sapph laser beam needs to be mode matched with the cavity. These matters will be considered sequentially in this section.
5.5. THE CAVITY

5.5.1 The Cavity Type

For several reasons, resonance of the cavity with both fundamental and second harmonic frequencies is highly desirable. According to ([74], eq.11) resonance of the latter would enhance the doubling efficiency by a factor

\[
\frac{1 - R_0}{1 - \sqrt{1 - L^2}}
\]

(5.16)

where \( R_0 \) is the reflectivity of the outcoupling mirror for the second harmonic and \( L \) is the loss per round trip for this wavelength. Another advantage of the doubly resonant setup is the TEM\(_{00}\) mode of the second harmonic beam, in contrast with the astigmatic beam generated in a singly resonant setup. However, a cavity resonant at the fundamental frequency is not necessarily resonant at the second harmonic frequency. Due to the difference in refractive indices in air for the two frequencies, one can achieve simultaneous resonance by varying the length of the cavity or by rotating the crystal. Small fluctuations in air temperature and pressure lead to the necessity to control two parameters in the doubling resonant cavity at the same time, e.g., path length (by translating one mirror) and orientation of the crystal, instead of only one as in the case of a single resonant frequency. As the adjustment of one of the parameters causes the need to re-adjust the other, the system is very complex to setup. Given this complexity, we have decided to use a singly resonant build up cavity.

Both standing wave and running wave cavities have advantages and disadvantages. Both configurations have similar intracavity doubling efficiencies. References ([63, 64, 73]) use a running wave cavity, and references [75, 74, 76] use a standing wave cavity. Several properties make a running wave cavity more attractive. The second harmonic light generated in the standing wave cavity propagates in both forward and backward directions. It can be rather difficult to recombine these two beams, which makes the usable power output of the standing wave cavity lower than that of the running wave cavity. Additionally, in a running wave cavity a plane mirror can be used to couple the light in, making the alignment easier, because the exact position on the mirror is of no importance. The position of the plane mirror has a smaller effect on the beam waist of the cavity than the position of a curved mirror, which makes the adjustment of the beam waist less critical than in the cavity with only curved mirrors. Furthermore, tilted curved mirrors can be used to compensate for astigmatism from a Brewster cut crystal. Probably the most important reason to choose a running wave cavity, is that
a standing wave will produce a large amount of back scattered light which can make locking of the laser very difficult, without further measures to reduce this.

Consequently we decided to use a running wave cavity, non resonant with the second harmonic. A schematic diagram of the cavity used in our experiment is given in figure 5.3. It shows a traveling wave cavity employing two curved and two flat mirrors, with the zero astigmatism in the region between the focusing curved mirrors.

5.5.2 The Cavity Design

As shown in figure 5.3 mirrors $M_1$ and $M_2$ are curved and $M_3$ and $M_4$ are flat mirrors. $l_1$ is the distance between the curved mirrors $M_1-M_2$, with focal lengths $f_1$ and $f_2$. $l_3$ is the distance between two flat mirrors $M_3-M_4$, and $l_2$ and $l_4$ are the distances between $M_2-M_3$ and $M_1-M_4$ respectively. We define: $L = l_2 + l_3 + l_4$. The stability parameter $G$ can be calculated [77] as

$$G = 1 + \frac{l_1 f_1 f_2}{2 f_1 f_2} - \frac{L + l_1 f_1 + f_2}{2 f_1 f_2}$$

and is independent of the beginning point of the round trip. For stability $|G| < 1$ is required, which gives the condition for the maximal $L$:

$$L < \frac{l_1 (f_1 + f_2)}{l_1 - f_1 - f_2}$$

Because of the tilted curved mirrors $M_1$ and $M_2$, the round trip matrices of the tangential plane (containing the chief ray and the optical axis and labelled $t$) and the sagittal plane (containing the chief ray, perpendicular to the tangential plane
and labelled $t$) are different. The focal lengths of a curved mirror (curvature $R$) tilted by $\Theta$ off the optical axis in the two planes are (eq 1, [78])

$$f_t = \frac{R}{2} \cos \Theta$$  \hspace{1cm} (5.19)

$$f_s = \frac{R}{2 \cos \Theta}$$  \hspace{1cm} (5.20)

The cavity is completely symmetric, so $f_1 = f_2$ for the same mirror curvature. According to ([78], eq. 33) the spot size of the cavity mode in the reference plane (beginning of the roundtrip) is given by the expression

$$\frac{\pi w^2}{\lambda} = \frac{2|B|}{\sqrt{4 - (A + D)^2}}$$  \hspace{1cm} (5.21)

The size of the waist (minimum spot size $w_0$) is given by (eq. 34, [78])

$$\frac{\pi w_0^2}{\lambda} = \frac{\sqrt{4 - (A + D)^2}}{2|C|}$$  \hspace{1cm} (5.22)

and the location of the waist relative to the reference plane by (eq. 35, [78])

$$z = \frac{A + D}{2C},$$  \hspace{1cm} (5.23)

where [77]:

$$A = 1 + x(\frac{XL}{f_1 f_2} - \frac{1}{f_1} - \frac{1}{f_2}) - \frac{L}{f_1}$$

$$B = L + x - \frac{XL}{f_2} + (l_1 - x)(1 - \frac{x}{f_1} - x f_2 - \frac{L}{f_1} + \frac{Lx}{f_1 f_2})$$

$$C = \frac{f_1 + f_2}{f_1 f_2} + \frac{L}{f_1 f_2}$$

$$D = 1 - L f_2 + (\frac{L}{f_1 f_2} - \frac{1}{f_1} - \frac{1}{f_2})(l_1 - x)$$

for a roundtrip beginning at a distance $x$ after mirror $M_1$.

The waist is therefore located midway between the two curved mirrors in both planes, independent of the wavelength. If the beam waists in both planes are the same, the cavity is non-astigmatic between the two curved mirrors, as the divergence of a Gaussian beam is completely determined by the waist and the wavelength. The condition

$$w_0^{(t)} = w_0^{(s)}$$  \hspace{1cm} (5.24)
is fulfilled if

\[ L = \frac{1}{2\alpha}(-\beta \pm \sqrt{\beta^2 - 4\alpha \gamma}), \]  

(5.25)

where:

\[ \alpha = 2cR - R(1 + c^2) \]  

(5.26)

\[ \beta = R(cR - l_1(1 + c^2)) \]  

(5.27)

\[ \gamma = cR^2 \]  

(5.28)

\[ c = \cos \Theta \]  

(5.29)

Therefore, \( L \) can always be chosen to make the astigmatism vanish between the curved mirrors. Since the solution to equation 5.2 does not depend on \( \lambda \), the cavity is nonastigmatic for any wavelength. One can easily obtain zero astigmatism in the other part of the cavity beginning with one curved mirror, leading over the two plane mirrors to the second curved mirror, by interchanging \( L \) and \( l_1 \) in the formulas.

The ray transfer matrix for a dielectric with refractive index \( n \) and length \( d \) and with surfaces perpendicular to the chief ray is given by ([77], table 1). In the case of type I critical phase matching, the fundamental has a single refractive index in a birefringent crystal, so the only effect of the doubling crystal on the \( ABCD \)-matrix of the cavity is a shortening of the length \( l_1 \) by \( l'_1 = l_1 - l(1 - 1/n) \). For an 18 mm long LBO crystal the beam waist providing maximum conversion efficiency is \( w_0 = 31.5 \mu m \) (see section 5.2).

This value for \( w_0 \) corresponds to the beam waist of a nonastigmatic cavity with the parameters

\[ l'_1 = 0.1069 \text{m} \]

\[ L = 0.7128 \text{m} \]

\[ \Theta = 0.113 \text{rad} \]

The sizes of the waists between the plane mirrors are \( w_s = 311.1 \mu m \) and \( w_t = 283.6 \mu m \).

### 5.6 The Locking Scheme

The idea behind the locking of the laser to the cavity, or vice versa, is to have the doubling cavity and the laser cavity coupled together to be resonant at the
same frequency. Knowing that the invar laser cavity is far more stable than the doubling cavity, we have locked the doubling cavity to the Ti-Sapphire laser.

The locking scheme used in our experiment was introduced by Hänsch and Couillaud [79]. The idea is to couple the two cavities by polarisation spectroscopy of the beam reflected off the input coupler. Using a relatively simple setup, without any modulation, this scheme generates a dispersive signal with sufficiently broad wings and very high signal to noise ratio. This means that it is possible to re-lock quite easily, even after larger accidental frequency jumps. The fact that only the beam reflected off the input coupler mirror is used for the locking means that the incoming beam, and consequently the second harmonic beam, are unaffected by the locking.

A schematic diagram of the locking scheme is given in figure 5.4. In theory, there should be a polarisation dependent element in the cavity (a Brewster cut crystal in the Hänsch-Couillaud scheme, or a thin Brewster plate in the Stony Brook group's experiment). Although we did not have a Brewster cut crystal, we had enough dispersion from the mirrors without introducing a new dispersive element in the cavity.

Consider linearly polarised light incident on the tilted input coupler, as illustrated in figure 5.4. The component with parallel polarisation to the mirror surface sees a cavity of low loss and experiences a phase shift due to the multiple reflections in the cavity. The component perpendicular to this polarisation simply gets reflected. The reflected beam now consists of two perpendicular polarisations

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**Figure 5.4:** Schematic diagram of a frequency doubler apparatus and the locking scheme
CHAPTER 5. GENERATION OF 389 NM LIGHT

with a relative phase difference. If the cavity is off resonance, the phase difference is non-zero and the two polarisations add up to give an elliptical polarisation. The handedness of the ellipticity depends on the sign of the detuning from resonance. To detect the ellipticity, the light is sent into the arrangement shown in figure 5.4. The fast axis of the λ/4 plate is rotated by 45° relative to the polarisation axis of one of the beam splitter outputs. The elliptical polarisation can be considered as a superposition of two counter-rotating circular polarisations of different amplitudes. The λ/4 plate transforms these circular components into orthogonal linearly polarised waves, separated by a beamsplitter and then monitored by two photodiodes (P.D.). Thus a linear polarisation consisting of two circular components with the same amplitude will generate two linear polarised waves with the same amplitude. This can be used to balance two photodiodes. The difference signal is zero for cavity resonance and the sign shows the sign of the detuning from resonance. The signal from the two photodiodes is subtracted, amplified and fed back to a piezo stack on one of the mirrors via a P.I.D.

5.7 Implementation

The experimental setup is shown in figure 5.4. Argon Ion Laser (Coherent Inova 300) was used as a pump laser for the Ti:Sapphire Laser (Coherent 899). The Ti:Sapphire output, of approximately 1 W single frequency, is mode matched by a spherical lens into the enhancement cavity; the polarisation of the input beam is controlled by a λ/2 plate. The nonlinear crystal is an 18mm long LBO crystal (AR coated for the fundamental and the second harmonic). The optimal cavity waist according to section 6.2 is:

\[ w_0 = 31.5\mu m \]  

(5.30)

for optimal focusing. This will result in a second harmonic coefficient (see section 5.2)

\[ \gamma_{sh} = 1.52 \cdot 10^{-4} W^{-1}. \]  

(5.31)

The circular waist is produced by focusing the beam using two curved mirrors with curvature radius 100 mm in a running wave bow-tie resonator \( (l_1 = 0.1069 \, m, \, L = 0.7128 \, m, \, \Theta = 0.113 \, \text{rad}) \). The blue light produced is then transmitted by the first mirror M1 after the doubling crystal, as this mirror has a dielectric HR coating for the fundamental light only.
The doubling cavity should be as stable as possible. Therefore the laser beam is guided close to the optical table, on specially shortened posts and holders, produced for us by the local workshop. In order to achieve phase matching by angle tuning, the crystal has to be rotatable about the vertical axis. Using a five axis manipulator (New Focus), the crystal can be translated in all three directions and rotated about the vertical axis and the axis perpendicular to the vertical and optical axis. Because of potential thermal problems with LBO in the regime of tight focusing [64], the crystal is mounted on a home built copper block as a heat sink using indium foil to increase the heat conductance.

First, the length of the cavity (without the doubling crystal) is scanned using a ramped high voltage on the piezo, and the transmitted fundamental intensity through mirror $M_1$ is observed using a photodiode. A lens is used to ensure that all transverse modes reach the photodiode. Although the reflectivity of mirror $M_1$ for the fundamental is very high ($r > 99.7\%$), an attenuator is found to be necessary to ensure the photodiode operates in a linear regime. The cavity length and alignment are now optimised to yield the largest intensity of the transmitted light on the TEM$_{00}$ resonance, and to suppress other transverse modes. As expected the finesse of the cavity is very high, producing a sharp resonance.

With the crystal in the cavity, pulsed blue light is observed as the cavity was scanned through resonance. In figure 5.5(a) the intensities of the transmitted fundamental light (red) and the second harmonic light (blue) are shown as a function of the piezo voltage. The smaller peaks in the fundamental transmission curve correspond to higher order transverse modes. Since these modes cannot be focused as tightly, even at equal powers they would have much less intensity. As the second harmonic power generated is proportional to the square of the fundamental intensity, the higher transverse modes are not visible in the second harmonic curve.

The finesse of the cavity can be calculated from the reflectivities of the cavity mirrors as:

$$\mathcal{F} = \pi\sqrt{r/(1 - r)},$$

with $r$ being the geometric mean of the reflectivities [77]. From the ratio of the FWHM of the Lorentzian fundamental transmission peak (see figure 5.5(b)) to the distance between two peaks we determine the finesse to be $\mathcal{F} = 120$. This corresponds to $r = 0.975$, which is in reasonable agreement with the input coupler reflectivity of 98%. The shape of the blue light transmission spectrum
Figure 5.5: The transmission intensities of the fundamental (red) and the SHG (blue) as a function of the voltage on the piezo. The finesse for the fundamental is \( \sim 120 \). The intensity calibration for the two traces is arbitrary, and the traces are offset. The line in (b) represents a fit to a standard Lorentzian (fundamental) and to the square of the same Lorentzian (SHG).
is exactly the square of the fundamental transmission spectrum, as illustrated in figure 5.5(b).

Slight vibrations and fluctuations can cause a cavity to be off resonance, which means that the cavity has to be stabilised. To increase the passive stability of the setup, the cavity was placed in a perspex housing, eliminating air currents and temperature fluctuations.

Due to the difference in characteristics of the first cavity eigenmode and Ti-Sapphire output, the input coupling is not very efficient. In order to minimise the reflection of the Ti-Sapphire output and maximize the input coupling, some mode matching is required. This requires the transformation of the Ti-Sapphire Gaussian beam into the cavity mode. We measured the waist of the Ti:Sapphire and simulating the cavity with a beam waist obtained from measurements, found that the optimal position of the mode-matching lens would be at 70 cm from the waist of the cavity. The lens was AR coated for 800 nm. Measurement with the pinhole on a translation stage with μm precision showed following results for the vertical and horizontal beam waist:

\[ w_{0v} = 420 \mu m. \]
\[ w_{0h} = 405 \mu m. \]

at distances of \( b_v = -0.443 \) m and \( b_h = -0.712 \) m from the laser output coupler, where \( w_0 \) values are \( 1/e^2 \) radii values.

In general, one spherical lens is not sufficient to match two astigmatic beams to each other. In order to further increase efficiency a spherical lens should have been used to match one plane and to correct the other plane a cylindrical lens should be used.

The cavity is locked to the laser frequency using the scheme described in section 5.6. The polarisation of the incident light is controlled by a \( \lambda/2 \) plate. In the current setup, the fundamental light is polarised in the tangential plane. By adjusting the \( \lambda/4 \) plate in figure 5.4, a symmetric error signal can be obtained, as in figure 5.6. The zero crossing generated is very sharp, and a tight lock should be possible. Again, the smaller features in the curve are due to higher order transverse modes.

The ramp on the piezo is now stopped and the cavity is locked to the fundamental laser frequency by feeding the error signal via a P.I.D. circuit and an amplifier to the piezo. The average power in the 389 nm output beam is now typically \( \sim 30 \) mW for an input power of \( \sim 900 \) mW. The output beam is elliptical.
CHAPTER 5. GENERATION OF 389 NM LIGHT

Figure 5.6: The error signal as a function of the piezo voltage (purple). The fundamental intensity is shown in red for reference.

and astigmatic. There is a fair amount of intensity noise (~ 30%) on the output beam, resulting from the still imperfect lock of the cavity to the fundamental laser frequency and from intensity noise on the fundamental Ti:Sapphire light. The 389 nm output radiation can now easily be tuned by tuning the fundamental laser frequency. Continuous frequency scans of the 389 nm light of up to 1 GHz, limited by the adjustment range of the piezo, are thus possible. Such frequency scans are used in the experiments in chapter 6.

5.8 Conclusions

We have obtained an output power of ~30 mW of 389 nm light, using a Ti:Sapphire laser which is frequency doubled in an external build up cavity. The power achieved is sufficient for the planned experiments, as we have very slow atoms (100 ms⁻¹) and consequently long interaction times. We have then locked the doubling cavity to the Ti:Sapphire laser applying locking scheme described in section 5.6, using a home-built PID amplifier and photodetectors. We have used the 389 nm radiation to modify ultracold collisions of cold helium atoms in a magneto optic trap, as detailed in chapter 6.

The conversion efficiency is significantly lower than ultimately expected. One of the key points for future enhancement of the doubling cavity, if more power is needed, is to improve the mode matching between the Ti-Sapphire laser beam
and the doubling cavity. This can be achieved by using a combination of cylindrical and spherical lenses and/or astigmatism compensated rhombs for the mode matching.

Better intensity stability of the output may be obtained by intensity stabilisation of the Ti:Sapphire laser, and improved locking of the build-up cavity to the Ti-Sapphire laser. The former can be implemented using a fast laser power detection circuit, feeding back to an acousto-optic modulator that deflects a controlled fraction of the incident beam. The latter can be achieved in two ways:

1. Use an amplifier with a larger bandwidth in the lock loop. Present bandwidth is \( \sim 8 \text{kHz} \), limited by the piezo capacity. A higher bandwidth would result in better noise suppression.

2. Reduce the amount of noise that needs to be regulated. This can again be achieved in two ways:
   
   (a) More stable mirror mounts in the build-up cavity

   (b) A more stable Ti:Sapphire laser. This could be achieved by using an EOM in the Ti:Sapphire cavity to provide high-bandwidth error correction, and an external high finesse etalon to provide high-bandwidth detection. A simple solution may be to use a fibre to couple light to the reference etalon. In this way, jitter in the pointing stability of the laser is not directly translated to frequency jitter by the locking electronics.
Chapter 6

Trap experiments

6.1 Introduction

Since the first magneto-optical trap (MOT) was achieved in 1987 [80], many groups have exploited the relatively high density ($10^{10}$–$10^{12}$ cm$^{-3}$) and low temperatures ($\sim 10\mu$K) of the trapped atoms. A large number of experiments have been carried out that utilise these traps as a source of extremely cold atoms. A review of the techniques used to achieve high densities and low temperatures in alkali MOT's can be found in reference [81]. We have developed a MOT for metastable He($^2S$) which is loaded by atoms from the bright beam source. The 389 nm laser source constructed in the course of this work (see chapter 5) has been used to study the trap and collision processes between the trapped atoms. A schematic representation of our magneto-optic trap is shown in figure 6.1.

Trap experiments with metastable atoms have demonstrated the additional problem of significant limitations to the trap density, due to ionising, and other inelastic, intra-trap collisions. These collision processes are an interesting example of the dynamics that may occur in ultracold samples of atoms in the presence of light. Several processes, such as Penning ionisation (P.I.) and associative ionisation (A.I.), have been identified:

$$\text{He}(^3S_1) + \text{He}(^3S_1) \rightarrow \text{He} + \text{He}^+ + e^-(\text{metastable P.I.})$$
$$\text{He}(^3S_1) + \text{He}(^3P_2) \rightarrow \text{He} + \text{He}^+ + e^-(\text{excited P.I.})$$
$$\text{He}(^3S_1) + \text{He}(^3P_2) \rightarrow \text{He}_2^+ + e^-(\text{A.I.})$$

The cross-section for these collision processes is very large: of the order of $10^6 \text{Å}^2$
6.1. INTROD

Figure 6.1: Schematic representation of a magneto-optical trap

for excited state Penning ionisation (PI) of cold atomic species. Hence these are dominant loss processes for atoms in the trap. This large excited state cross section PI can be explained using the knowledge (see figure 6.2) that the interaction in the quasimolecule, consisting of a metastable ($^{23}\text{S}$) atom and an excited state ($^{23}\text{P}$) atom is governed by a resonant $1/R^3$ potential, instead of the $1/R^6$ potential between two metastable state atoms. For negatively (red) detuned light, as is used to form the trap, the quasimolecule can be excited to an attractive $1/R^3$ potential. The atoms than accelerate towards each other, and may or may not decay spontaneously. Penning and/or associative ionisation can than occur at small internuclear distance. If the quasimolecule decays during the interaction, the main effect is that the collision energy is higher, and hence the number of partial waves contributing to the Penning ionisation process is larger. This is observed as an enhanced ionisation cross-section. If the quasimolecule does not decay during the collision, the possibility for associative ionisation exists [82]. Also, for this process, the collision energy has been increased by acceleration in the attractive potential.

On the other hand, for positive (blue) detuning the quasimolecule may be excited to a repulsive potential, thus reducing the ionisation rate. This has been demonstrated in an experiment done at NIST [58] on Xe. They have shown that probing the trapped metastable Xe atoms with blue detuned light decreases the
collisional ionisation rate by a factor of 8, when the trapping lasers are off, and a factor of 30 when the trapping lasers are present. A similar reduction may occur in metastable He.

Two prior experiments have studied the effects of near resonant light on trapped metastable Kr [19] and He [83] atoms. The latter experiment, analyses the trap losses by monitoring the ion current due to Penning collisions either with background gasses, or between trapped atoms.

The results with Kr have also confirmed the crucial role of the near resonant light on collision dynamics. The ionisation rate was shown to increase several times in the presence of the laser, red detuned by a few times the natural width, while it sharply decreases for a slight blue detuning.

The atom optics group at the Free University of Amsterdam has used 389 nm light to probe ultracold collisions of metastable helium atoms in a magneto-optical trap [84]. They have detected ions formed in the trap, while scanning the probe laser. It has been shown that the ion formation is significantly enhanced when the laser is close to the resonance for one of the $3^3P_{0,1,2}$ fine structure states (see

\[ E/h \begin{align*}
&\begin{array}{c}
\text{repulsive} \\
\text{attractive} \\
\text{attractive}
\end{array}
\end{align*}
\]

**Figure 6.2:** Potential energy curves for the $\text{He}_2^*$ quasimolecule when one of the atoms is excited to the $^3P$ state in an S-wave cold collision. The different curves represent different relative orientations of the atoms. At small internuclear distance, ($d < 100a_0$) the potential curve intersects with an ionic potential curve. The distance scale is in atomic units.
More recently, new experiments on the modification of ultracold metastable He collisions have been performed [18]. They confirm a significant increase of the collision rate for negative detuning. In another experiment at the Free University in Amsterdam the He atoms have been spin polarised using a circularly polarised laser. Theoretical calculations expect the ionisation rate to decrease by five orders of magnitude for spin polarised atoms [85, 86]. The Amsterdam experiments so far confirmed a supression by two orders of magnitude, limited by their signal levels. This supression may open up the possibility for the observation of a Bose-Einstein Condensation in metastable helium. However, the spin polarisation can only be performed once the atoms are released from the MOT, and cannot be used to obtain a higher trap density in the MOT itself.

Our experiment is aimed as an additional study of the collision dynamics in
the presence of a catalysis laser, tuned to the $^{2}S \to 3^{2}P$ transition at 389 nm. We
trap the metastable He atoms using the infrared 1083 nm transition. However,
we extinguish the trapping lasers during the experimental period, enabling a
much clearer comparison with theory, as was done in the NIST experiment on
metastable xenon [58]. As a side experiment we have repeated the experiment of
the Amsterdam group in which the probing was done in the presence of trapping
fields [20].

6.2 The design of the trap

The trap is designed to work as a “classical” six-beam Magneto-Optical Trap.
The geometry of the setup is shown in figure 6.1, and the hardware of the trap,
 machined in our workshop, is shown in figure 6.4. The distance between the end
of the slower and the center of the trap is 270 mm. A pneumatic valve between the
trap and the Zeeman slower made it impossible to have the trap any closer to the
end of the slower but, on the other hand, it allowed more flexibility in adjustments
of the apparatus during the experiment. The trap is continuously loaded by the
slowed atomic beam leaving the decelerator. The atoms are trapped using three
orthogonal pairs of circularly polarised, counterpropagating laser beams. The
geometrical position and orientation of the fields employed in the trap is shown
in figure 6.1. Besides the quadrupole trapping magnetic field generated by two
anti-Helmholtz coils, there is a magnetic field coil opposing the residual field of
the Zeeman slower to give a zero field in the z direction in the middle of the trap,
as well as three pairs of anti-Helmholtz coils for fine tuning of the zero field in
the middle of the trap. The anti-Helmholtz coils for the trapping field are at a
mutual distance of 220 mm, formed by 250 windings of wire with a cross-section
of 1.6 mm on a 74 mm diameter former, which is water cooled. The typical
current through each of the coils is 5 A, and the average diameter of the coils
is 100 mm. Whilst this is far from the true anti-Helmholtz configuration, which
would yield the strongest field gradients, the current setup is a compromise which
enables excellent access to the trap for the many laser beams.

With the trap in place the 'slowing mirror', coupling in the laser beam for
the Zeeman slower, is placed downstream of the trap and slowing is done by a
separate laser. The trapping beams are generated by a single SDL diode laser,
with a total power of 50 mW. The axis of the two diagonal (vertical) trapping
6.3. OBTAINING THE MOT

Figure 6.4: The trap chamber

beams are at an angle of $45^\circ$ with respect to the atomic beam, and the horizontal beam is perpendicular to the atomic beam. All three beams are retroreflected by mirrors positioned outside the trap chamber with quarter wave plates in front of the retroreflecting mirrors. The diameter of the beams is around 2 cm, with an average intensity of 1.6 mW cm$^{-2}$ or $I = 10I_0$. Both the trapping and probing beams are switched using AOMs for which the drivers have been developed during this project see (section A.3)

The MOT chamber is a high vacuum chamber pumped by a turbo pump (Pfeiffer TPH-180), backed with a CIT Alcatel turbo pump. The typical working pressure of the chamber in the presence of the atomic beam is $2 \times 10^{-8}$, whilst the base pressure is $5 \times 10^{-10}$ mbar. The windows of the MOT chamber are antireflection coated for the 1083 nm laser wavelength.

6.3 Obtaining the MOT

The 1083 nm light we use to cool and trap the atoms can be detected by a standard CCD camera, albeit with a very low quantum efficiency. Hence, observing infrared 1083 nm fluorescence is almost impossible for a partially working trap. To be able to optimise the trap we illuminate the center of the trap region with 587.8 nm light, generated by a dye laser, and resonant with a visible $2^3P_2 \rightarrow 3^3D_J$ transition in He. As the lower state of this transition is the $2^3P_2$ state, the trap lasers are
required to be present as well. Initially, the yellow laser is tuned to resonance by observing fluorescence in the collimated, but not slowed atomic beam using an intense yellow laser beam. Then the slowing laser is turned on, the intensity of the yellow beam is decreased, and we tune the field of the second half of the Zeeman slower to obtain a MOT. This initial MOT is very weak, due to the non-perfect overlap of the laser beams and residual magnetic fields. Once the initial MOT is found, it can be optimised using the Helmholtz compensation coils and by adjusting the laser beam overlap. As soon as the trap has been optimised in this manner, the yellow laser is no longer required as the 1083 nm fluorescence becomes visible. The trap is very reproducible from day to day.

Once the trap contains a sufficient number of atoms, it can be visualised using either the 588 nm, 1083 nm or 389 nm transitions. Such images of a well-working trap are given in figure 6.5. As the detection efficiency at 1083 nm is still quite poor, there is significant noise in this image.

![Images of a trap fluorescing under different transitions.](image)

*Figure 6.5: The trap fluorescing under excitation of different transitions. All images were recorded in black and white with a standard CCD camera and have been false-coloured.*

The force confining the atoms in the trap is the largest for a high laser intensity and in the presence of high magnetic field gradients. Since the transverse spreading of the atoms coming out of the Zeeman slower is significant, the large laser beam diameters are used in order to make the geometric capture range of the trap as large as possible. The first images are taken with the 50 mW diode laser power equally divided between the vertical and horizontal beams, a laser
detuning of 42 MHz, a laser beam diameter of 2 cm and a magnetic field with a gradient of 40 Gauss/cm. The alignment of the lasers, and their polarisation, have proved to be noncritical. This detuning, equal to ~25 linewidths, is extremely large compared to traps for alkali atoms. The large detuning yields a large diameter, low density trap. Still, these parameters seem to yield the largest number of atoms, as determined from the fluorescence. This can be explained by the very large cross-section for ionising collisions, particularly if one of the collision partners is excited by the laser. Indeed, we have found that the ion count from a channeltron on the bottom of the trap chamber is a much more sensitive tool to detect the presence of trapped atoms than the fluorescence from these atoms on the 1083 nm transition.

6.4 Cold Collisions Experiments

As mentioned in the introduction, the trapped atoms are expected to undergo Penning- and associative ionisation processes. The laser fields used for trapping can strongly modify these processes. If a second laser field is present, which excites a different electronic transition, and which has a different detuning from this transition, the analysis of a collision process becomes very complex. Hence, we have chosen to do measurements where, if possible, only one laser field (either probe or trap) is present. In this work, we describe the modification of the ionising collision rate due to the presence of the 389 nm laser (see chapter 5), tuned close to resonance with the $2^3S \rightarrow 3^3P$ transition.

The laser is tuned to the resonance line in three steps:

1. The frequency of the Ti:Sapphire laser is measured using a home-built wave meter [87], and tuned to half the known frequency of the transition [88]

2. The 389nm laser beam is sent through a helium discharge absorption cell. As the absorption is too small to measure, the fluorescence is monitored through a monochromator, tuned to the 389 nm fluorescence line, as the laser frequency is scanned. The fluorescence increases strongly as the laser scans through the Doppler broadened absorption profile of the transition in the cell.

3. Observation of the fluorescence in the trap.
As mentioned, the absorption in the cell is very small, and the 389 nm laser has significant intensity noise. This made it impossible to determine the line centre with any accuracy, or to use saturated absorption spectroscopy to find the line centre and calibrate the detuning.

Within the Doppler profile of the emission from the cell, the frequency of the probe laser, as determined by the frequency of the fundamental Ti:Sapphire laser, can be changed by putting a voltage on the external input of the Ti:Sapphire controller box. The doubling cavity is locked to the fundamental laser frequency and the frequency calibration is obtained from the frequency scale of the Ti:Sapphire laser. The detuning of the probe laser is thus scanned over a range of 500 MHz around the atomic resonance, including the resonance. Currently, we have no method to determine the absolute frequency of the resonance. The probe laser power is typically in the range of 10–30 mW, with a beam diameter of 3 mm. The probe laser is retroreflected to both equalise the light forces on the trapped atoms (minimise the disturbance of the trap) and to increase the probe laser intensity.

The first collision experiments in the trap were monitored using a channeltron (1 cm in diameter), placed at the bottom of the trap chamber. The voltage on the entrance of the channeltron is -300 V, and +2700 V is applied at the rear of the channeltron. In order to selectively detect metastable atoms we used a grid in front of the channeltron, at a potential of 20 V.
To determine the modification in ion production rate by a probe laser, we have used the following scheme, which is illustrated in figure 6.6:

1. The 389 nm probe laser is chopped using a mechanical chopper, creating nearly identical 'laser on' and 'laser off' times. A small part of the probe laser beam is split off, and hits a photo-diode. This timing signal is used as a basis for all other time periods.

2. The trapping lasers are turned off during the experiment period (100 μs) using an acousto-optic modulator

3. Two ion count signals are obtained, one with the probe laser off (period $A$) and one with the probe laser on (period $B$). The experiment period is sufficiently short so that the atom density does not change significantly during this period.

In the experiment, the laser frequency is stepped, and two separate GPIB counters are started by a computer program. Counts are accumulated during a period of one to four seconds, representing several thousand laser on–laser off cycles. Longer data accumulation times led to unacceptable laser drift during the experiment. Subsequently, the number of ions detected is read out of the counters by the GPIB interface board in the computer. This procedure is repeated for each probe laser detuning.

### 6.5 Cold Collision Results and Discussion

The ratio of the number of counts in the presence of the probing laser and in its absence, giving us the proportion of metastables in respect to the background ions, is given in figure 6.7.

The ions detected during period $A$ can be due to collisions of metastable $^3S$ atoms either with each other or with background gases. With the current equipment we cannot distinguish these two processes. Experiments performed elsewhere [18] suggest that a large fraction of these counts are formed by the ionisation of background water vapour. However, the difference between the ion counts during the periods $A$ and $B$ is only due to the presence of the probe laser. As the mechanism for strong enhancement of the ionisation rate, as suggested in the introduction of this chapter, is only valid for $^3S \rightarrow ^3P$ collisions, we presume that the probe laser has a negligible effect on the ionisation of background gases.
CHAPTER 6. TRAP EXPERIMENTS

Figure 6.7: Ratio of laser-on and laser-off ion count rates under different conditions

(a) Different magnetic fields in the trap

(b) Different probe laser powers (intensities)

(c) Different trapping laser field detunings
The ion count rate shows a strong enhancement around resonance, up to a factor of two. With the current frequency reference setup it was impossible to determine whether the maximum in the count rate coincides with the resonance frequency, or that the maximum is slightly below resonance, as expected. In the graphs, the frequency offset on the x-axis is adjusted such that the count-rate maximum occurs at the same frequency. Significant adjustments were needed due to drift of the Ti:Sapphire laser frequency.

In figure 6.7(a) the ion count rate as a function of the probe laser frequency, at a number of values for the probe laser power, is displayed. We can observe that the width and the height of the peak increase when the probe laser power is increased. This can be understood qualitatively from the following argument: In a dressed atom picture we label the $2^3S$ asymptotic state as $|1\rangle$, the $3^3P$ state as $|2\rangle$, and the number of 389 nm photons as $|n\rangle$. The crossing of the $|1,1,n\rangle$ state with the $11,2,77. - 1\rangle$ state changes into an anticrossing as the coupling between these is increased. The latter state is now on the attractive $-1/R^3$ potential that gives rise to the increased ionisation rate. However, the low signal to noise ratio prohibits any more quantitative statements at this moment.

In figure 6.7(b) we show the ion count rates as a function of probe laser frequency for different magnetic field gradients. This should result in a different atom density in the trap, although no density measurements were performed to verify this. No significant difference is observed between the curves.

Figure 6.7(c) shows several curves for different trap laser detunings. A larger trap laser detuning leads to a higher temperature and lower density of the trapped atoms. Also, the absolute number of atoms, as determined from the brightness of the fluorescence from the trapped atoms, goes through a maximum for $\Delta \approx 26$ MHz. The main difference between the curves is a different ratio height at the peak, reflecting the different number of trapped atoms. Again, further measurements with a better signal to noise are required for any quantitative conclusions.

In none of the curves do we observe suppression of the ionisation rate (expected for blue detuning). This is in agreement with experiments done elsewhere, but at variance with the simple physical collision picture presented in section 6.1. This can have number of causes:

- The laser off ion rate could contain a large component of background molecular ions, and these collisions are not suppressed. In that case, the signal
to noise ratio in the experiments is currently insufficient to observe changes to the small fraction of the ion count rate that is due to He-He collisions.

- The long lifetime of the excited $3^3P$ state (100 ns) and the relatively high temperature of the atoms yield significantly different collision dynamics, that can not be explained by the simple model.

The signal to noise is currently too small to rule out either of these explanations. However, the first hypothesis is the simplest, and hence seems currently the most likely explanation. Future experiments could use either time-of-flight or quadrupole mass spectrometry to discriminate between the different ion products, and this way make more definitive statements.

6.6 Conclusions and suggested improvements

In summary, we have built a magneto-optical trap for metastable helium atoms that can be loaded from the "Bright Beam" discussed in chapter 5. For the current conditions, the trap contains $\sim 3 \cdot 10^6$ atoms in a 3 mm diameter cloud, at a temperature of $\sim 1$ mK. The trap density was considered to be limited by intra-trap collisions, particularly collisions between one atom in the $3^3P_2$ state and one in the metastable $3^3S_1$ state. As the former atoms are continuously created by the lasers we use for the trapping mechanism, this is hard to avoid. A tentative solution was found in using a large laser detuning in the trap. Hence the excited state fraction of the atoms is small, and the collisional loss is reduced. A limit to this technique is formed by the size of the trapping laser beams. This also means that the trapped cloud is relatively dilute compared to magneto-optical traps for alkali atoms. I suggest the following improvements to obtain a more dense and stable trap:

1. Increase the incoming flux of atoms by focussing and/or compressing the incoming atomic beam

2. Use a shutter to cut off the gas flow during the experiment phase: this should decrease the collision rate with background He gas

3. Use larger laser beams to increase the capture range of the trapping process

4. Use optical feedback to further stabilise the laser frequency of the trapping lasers.
Preliminary experiments have been carried out to study the effects of an exciting laser to the ionising collisions in the trap. We have used the 389 nm laser that has been developed to excite the $2^3S \rightarrow 3^3P$ transitions in helium. The laser is shown to increase the ion production rate by up to a factor of two. A decrease of the ionisation rate, as expected for a blue detuned laser, was not observed. The low signal levels during these experiments have so far limited any more quantitative observations. I recommend the following steps to improve these experiments:

1. Better intensity stability of the 389 nm laser light, as detailed in 5.

2. Improved frequency stability of the 389 nm laser. This is directly related to the previous point. Additionally, an absorption cell with significant absorption on the 389 nm line could provide the required frequency calibration.

3. An overhaul of the ion detection system to increase the count rates and enable time-of-flight mass discrimination for the detected species.

4. A shorter duty cycle for the 389 nm light to avoid disturbance of the trap by this laser, consequently this would enable more detailed measurements closer to resonance.

5. A more stable trap as detailed above.
Chapter 7

Conclusion

In this thesis I report on the successful construction of a bright source of slow metastable \( (2^3S) \) helium atoms. Laser Doppler cooling is used to increase the atomic beam brightness in three stages: the initially diverging beam from a LN\(_2\) cooled discharge source is collimated using effectively curved wavefronts. The resulting collimated beam, which has captured a solid angle of 0.005 sr is subsequently slowed in a Zeeman slower to \(~100\) m/s. The emphasis in this thesis lies with the third stage, in which the slow beam is compressed in a two-dimensional magneto-optic trap, and the collimation of the beam is improved using a two-dimensional optical molasses. The beam flux achieved is in excess of \(10^{10} \text{s}^{-1}\) in an area of \(~10\) mm\(^2\), whilst the divergence is a few mrad. This beam provides an excellent source for a range of “user” experiments, some of which are already underway.

In a preliminary experiment, I report on the study of one dimensional polarisation gradient dependent laser cooling of sodium atoms. The transverse velocity distribution of polarisation gradient cooled sodium atoms at \( \{3s\}^2S_{1/2} \ (F=2) \rightarrow \{3p\}^2P_{3/2} \ (F=3) \) transition at 589.1 nm was measured for a range of laser fields. Detailed analysis of experimental results gave us a better insight into dependence of efficiency of the polarisation gradient cooling process in regard to cooling laser intensity, detuning and polarisation. Sub-recoil resolution measurements have been performed on the transverse cooling of a longitudinally velocity selected sodium atomic beam in a lin\(_\perp\)lin laser field. The insights acquired in these experiments have been used in the design of the post-collimation, compression stage of our bright beam source, to further narrow the transverse spread of the already compressed, slow beam of helium atoms.
7.1. FUTURE RESEARCH

In parallel with the development of a bright beam machine a laser source for 389 nm UV light, which corresponds to $2^3S_1 \rightarrow 3^3P$ transition in helium was developed. The fact that this transition has a recoil velocity which is three times bigger recoil than that of the 1083 nm transition ($25.6 \text{ cm/s}$, $v_{rec}/v_{Dop} = 0.7$) was one of the reasons for this development. The thesis reports on frequency doubling of the output of a 778 nm Th-Sapphire laser beam, in an LBO crystal in an external build up cavity.

Finally, the thesis reports on the development of a Magneto-Optical-Trap (MOT) which is loaded by the collimating and cooling stages of the bright beam source. Our experiment was a study of the cold collision rates in the presence of catalyst laser. While trapping metastable He atoms using the infrared 1083 nm transition, we have studied the change of the ionising collision rate whilst probing with a different colour laser (389 nm) near resonant with the molecular potentials, in the absence of the trapping laser fields.

7.1 Future research

Both the "bright atomic beam" and the "metastable atom trap" form excellent sources for a range of "user" experiments. As an example, an experiment that is currently underway is the study of the guiding of metastable atoms in hollow optical fibers. Also, this year (1999) the first electron scattering experiments are expected to start. These experiments depend critically on the unique parameters of our bright beam source: The atomic beam has a high density, is well collimated and has a narrow longitudinal velocity distribution.

The results from the cold collision experiments in the trap are at present not well understood, and warrant an extensive follow-up. If the Penning ionisation rate can be sufficiently suppressed by the presence of additional laser fields, this will enable a large increase in trap density - of great importance towards the achievement of Bose Einstein Condensation in metastable helium. The trap will also be used as a source of ultracold atoms for atom optics experiments. The 389 nm light will play a key role in many of these experiments. The suggested improvements in section 5.8 on the 389 nm laser source should result in higher achievable laser power and power stability, and should greatly facilitate the employment of the 389 nm radiation in sensitive experiments.

The improvements on the "bright beam apparatus" suggested in section 4.5.6
could improve the beam flux and density even further. The electron scattering experiments depend strongly on the atomic density, and are thus anticipated to benefit most from these developments. However also in atom optics experiments, where better beam definition can be used without sacrificing signal levels, and atom lithography experiments are expected to benefit from a higher beam flux.
Appendix A

Electronic devices developed for the bright beam machine and related experiments

A.1 Diode laser stabilisation

Atom optics experiments necessitate a significant input in the design and development of electronic devices, such as stable and reliable drivers and locking systems for the lasers used in the experiment, as well as a number of smaller circuits, like detectors and PID (proportional-integral-differential) controllers. For some experiments it is also necessary to significantly shift the laser frequency, which is usually done by employing an AOM (acousto-optic modulator), which also requires driving electronics.

In the present studies several complementary methods for improving the stability of a free running diode laser were undertaken. We were primarily concerned with: (a) the construction of a stable power supply to reduce the mains frequency and RF noise, and (b) long term drift of the laser frequency, for which we have developed a lock-in amplifier to be combined with a simple PI circuit. While we have kept everything simple and inexpensive to construct, the stability is comparable with that of far more expensive, commercially available systems. Development and testing of our stabilisation scheme was done on the Distributed Bragg Reflector (DBR) diode lasers for use in a range of atomic collision and atom optics experiments, but it is applicable to all diode lasers, and can be used in a large range of atomic physics experiments.
We use helium metastable \((2^3S)\) atoms, which have a very long lifetime (theoretically predicted estimate 8000 s [89]), and which possess a number of convenient laser transitions for light force manipulation. We employ, primarily, the \(2^3S_1 \rightarrow 2^3P_2\) transition at 1083nm for which commercial distributed-Bragg-reflection (DBR) lasers (50mW) are available (SDL 6702-HI). The natural linewidth of the atomic transition is \(\Gamma/2\pi = 1.6\) MHz which can be resolved by careful stabilisation of the diode laser frequency.

### A.2 Diode laser controllers

The frequency of a free running diode laser shows high sensitivity to the operating temperature (in our case 18 MHz/mK) and to the injection current (with a typical response of 2 GHz/mA [90]). A number of current controllers on the market are able to provide the necessary current or temperature stability, but they require a perfectly “clean” power supply. These supplies, and consequently the whole control system, are generally very expensive. We use a design for a relatively inexpensive low-noise, high-speed laser current controller, developed by Libbrecht and Hall [91] together with a commercially available temperature controller LDT 5100 manufactured by ILX. Both of these are powered by commercial bench-top supplies. We have introduced some changes in the design of the current control board appropriate for our usage of the board. The schematic diagram of the current controller is given in figures A.1, A.2 and A.3. We have designed an interface board for the temperature and current controller which contains inputs for external modulation via buffer and integrator circuits, and front panel controls. A detailed circuit diagram of this board and interfaces to the temperature and current controllers, and to the front panel, are given in Figure A.4.

#### A.2.1 Circuit description

**Laser current control**

The current control boards (ILX LightWave LDX 3100), that were originally purchased to drive the laser diodes, turned out to be too noisy for our experiment. This initiated the development of a new current controller board.

The heart of the new diode current controllers is the high precision, constant current circuit designed by Libbrecht and Hall and modified to add protection
A.2. DIODE LASER CONTROLLERS

capability to the original design. This design relies on a very stable reference diode (U11) and a very low noise operational amplifier (U6) for its performance. Most of the remainder of the circuit is intended for monitoring (U8, U9A, U10A) or modulation (U3, U4, and U7). DC current feedback is provided by means of R38 which is made up of two 50 W, 5ppm/°C resistors in parallel. The laser current is converted to a voltage by R38 and matches the control voltage on the other input of U6. Modulation of the laser current is by means of U3, U4, and U7. This particular circuit looks quite odd at first glance but is designed to sum the current through R3 with the current provided by U6 and Q2.

The protection added to the original circuit was designed to emulate the protection used by the original ILX LightWave LDX 3100 laser current controller. The voltage across R38 is sensed by U8 and compared to a set point established by means of U5. Over current results in the analog multiplexer (U12) being switched to the current limit rather than the current set value. This limit is not latched. Open circuit on the other hand is latched by flip flops (U14) to disable the laser drive via Q7, Q8, U13, and D4. Because of the latching of the flip flops it is necessary to cycle the power to remove the overvoltage condition. Failure of either rail voltage will also disable the laser current by means of U9C and D as well as U13A and C and Q4.

The overvoltage set point is not adjustable and is set by R30 and R31. U9B compares the fixed set point with the output voltage of U6. Adjustment of the current maximum is by means of RV2 while RV3 sets the nominal laser current itself. The 3rd adjustment on the PCB is for the 10V supply generated by U1. RV1 is used to adjust the output of U1 to 10V as measured on TP3. In order to provide better cooling of the laser diode, some manufacturers electrically connect the anode and the housing of the diode. As current driver boards provide a positive current they are, as such, not directly compatible with this design. Our experience is that the safest construction is to have the housing (and thus the anode) floating. Hence, the diode laser housing is clamped on the heat sink using standard T03 electrical insulation materials. A properly designed heat sink still provides sufficient cooling for the diode laser.
Figure A.1: Diode current driver, sheet 1 of 3
Figure A.2: Diode current driver, sheet 2 of 3
Figure A.3: Diode current driver, sheet 3 of 3
Temperature control

The temperature control circuit is based on a LightWave LDT 5100 Temperature Controller. Its design will not be described here.

Laser current fine control

The feedback circuit is made using a number of operational (op) amplifiers (amps) and instrumentation amplifiers. The input signal is first buffered by instrumentation amp U1. This provides common mode rejection and impedance buffering. The output of this amplifier is then made available to both a differentiator stage (U2, U3) and to an integrator stage (U4, U5). The outputs of these two stages are summed by amplifier U6. Front panel pots allow adjustment of the gain of each section separately. The user can select to use a simple buffer (U7) by feeding signal into J3 (buffer input) on the front panel. Alternatively the signal can be fed to the differentiator/ integrator stages. The integrator stage has an adjustment for offset (VR2). The Fine Current pot on the front panel provides the same overall function by adding an adjustable offset to the output signal via U8.

Metering

The metering circuit is simply a 5V powered 3.5 digit LED digital panel meter which allows monitoring of either thermistor resistance (related to laser temperature) or the laser drive current. Since the meter module is 200mV full scale, VR4 and VR5, both on the front panel, are included to divide the input voltages down to that level. This allows trimming of both the temperature and the current measurements. A temperature/current toggle switch (Meter Sw) allows the user to select either the temperature or the laser current for monitoring.

Power supply

The main power supply for the system uses a Statronics Model 65/3 Linear power supply. The outputs of the supply are adjustable +V, -V and a fixed +5V. The adjustable supplies can be adjusted between 12 and 15 V. The mains connects via J1, a IEC connector with integral fuse, and is protected by a Slow Blow 500mA fuse. Two separate switches are used to turn the power on (at the request of the users). Front panel LEDs, D1 and D2, indicate that the +V rail is operating
Figure A.4: Feedback and control circuit
A.2. DIODE LASER CONTROLLERS

The mains (50 Hz), filtering scheme, as described in the text, uses constantly re-charged batteries, as a filtering element. The power supply is protected by the fuses.

correctly. Output connectors J2 and J3 are metal Cannon connectors which were chosen for shielding purposes.

A.2.2 Noise reduction and stabilised power supply

We have measured the noise of the current controller output using a spectrum analyser, and observed both RF and 50 Hz (mains frequency) components and their harmonics. To minimise this, we placed all our electronics modules in diecast boxes, and double shielded the cables from the control box to the laser housing. This employed coaxial cable plus metal braided shielding, which was extremely effective in reducing the RF noise.

A problem which has proved to be more persistent was 50Hz noise from the power supply. In principle there are two ways to solve this problem. One is to suppress the noise by adding filtering electronics, which for this frequency requires bulky inductors and capacitors. The other possibility is to run the controllers from large and expensive, rechargeable batteries. The Peltier cooler for the temperature controller would still need to be supplied from a 5V power supply, since it requires 1A of current. That leaves no more than 300mA required for the controlling electronics, but it still means that the batteries would have to be recharged every few days. So, we decided to make a simple and practical configuration (fig A.5) which allowed batteries to be constantly re-charged. We use small and inexpensive rechargeable batteries and similarly inexpensive benchtop power supplies. The idea is basically to keep the voltage of the power supply just below the maximum voltage on the terminals of the battery, so that the batteries
never get fully charged (in our case the batteries were charged up to 13.5V, and the maximum possible charge of the batteries is 13.8V) and act as a huge filtering capacitor, reducing the level of 50Hz noise on the diode laser current to the nA region.

**A.2.3 Lock-in amplifier**

To counter long term temperature drift, the laser frequency is locked to a reference, saturated absorption discharge cell [92]. An error signal is obtained by modulating the frequency of the transition using the Zeeman effect in an AC (~6 kHz) magnetic field, combined with phase sensitive detection. The error signal is fed back to the laser current control through a simple proportional-integrating (PI) circuit. We have designed a practical, uncomplicated and very effective lock-in system for the phase sensitive detection. A block diagram is given in figure A.6. It consists of a low gain amplifier for a reference signal (1), a phase-shifting section (2) after which the signal is converted to a square wave (3) and multiplied (4) by the photodiode signal, which in turn has been amplified by an adjustable gain amplifier (5). Both an amplified (7) and monitoring (8) output are provided.

A detailed schematic of the lock-in system is given in figure A.7, and it operates as follows. The (LF356N)-U1 is used both to provide low gain (×5) amplification of the reference signal, and to bring the voltage up, above 2V, for the input of the phase shifter. R3 and R4 are biasing resistors, and two low capacitance input FET's, Q1 and Q2, were used to form two 180 degree, phase shifters, with C2 as a DC separator in between. The phase shifter compares the reactance of
Figure A.7: Detailed schematic of the lock-in amplifier. For details see text.
Figure A.8: The result of a beating experiment performed with two SDL6720 diode lasers, with both lasers locked to the atomic transition using the method discussed in the text. The frequency range is 50 MHz, and the width of the beat signal observed is \( \sim 4 \text{ MHz} \).

C1 and C3 against the variable resistor VR1. Q3 is used as a driver FET and U2 as a square wave generator. C4 and R10 are used to bias the phase-shifted signal to ground to give a duty-cycle of 50% as the output of U2. The output signal from U3 is then multiplied by the error signal (arriving at J2). The error signal has been AC coupled (C5) and amplified through two amplification stages, for which two low noise op-amps have been used (OP-42). The first one, U3 is in a configuration which provides a gain of 3 whilst amplifier U4 has an adjustable gain from 1-500, depending on the value of the variable resistor VR2, resulting in an adjustable amplification from 3-1500. The multiplier is an inexpensive item (AD633). The output from the multiplier provides the input to the adder-amplifier U6 which adds an adjustable offset voltage from J3. U7 is employed in the configuration of a low pass filter-integrator with gain, with a time constant \( 1/RC = 2\pi \cdot 1 \text{ kHz} \). The resulting fast response is important to significantly decrease the diode laser bandwidth. The two outputs J4 and J5 represent the slow and fast outputs respectively. The capacitor network included in figure A.7 are the decoupling capacitors for the operational amplifiers.

A.2.4 Results

To measure the linewidth of the lasers with the above controllers we have performed a heterodyne beating experiment with two SDL-6720 lasers. One of the lasers is locked to the atomic resonance, the other to a detuning from resonance
A.3. AOM DRIVERS

of ~50 MHz. The 50 MHz beatnode is observed using a fast photodiode and an RF spectrum analyser. In figure A.8 the beat frequency spectrum over the range 25–75 MHz is displayed. The width of the observed Lorentzian peak is a direct measure of convolution of the linewidths of the two lasers. The FWHM of this peak is ~4 MHz, indicating a linewidth of a single laser of ~2 MHz, comparable to the linewidth (\(\Gamma/2\pi = 1.6\) MHz) of the \(2^{3}S_{1} \rightarrow 2^{3}P_{2}\) transition. This is to be contrasted to a linewidth of 30 MHz using the less stable current drivers, or 10 MHz using the stable current drivers without active locking.

A.3 AOM drivers

![Diagram of AOM drivers]

Figure A.9: Block schematic of the AOM drivers

The laser frequencies are very conveniently tuned using acousto-optic modulators (AOMs) each driven by a power amplified frequency source. The system design presented here allows the AOMs to be tuned to an arbitrary frequency between 100 and 150 MHz. The input power needed to drive this particular type of AOM (Crystal technology 3125-120) is 2 W. Synthesised or frequency locked sources, and voltage controlled oscillators (VCOs), are generally used to drive acousto-optic modulators. VCOs have significantly lower prices, and can provide the necessary frequency accuracy, which made them an obvious choice. Drivers for similar acousto-optic modulators with different tunability and input power requirements had already been developed and used in the Physics Department
Figure A.10: The dependence of the output frequency on the input voltage of the AOM drivers

(the Faculties) at the ANU [93]. We decided to adopt this design and improve parameters which we considered important for our experiment. Figure A.9 is a block diagram and figures A.11, A.12 and A.13 show detailed circuit diagrams of the AOM driver’s design.

The voltages used to control both the frequency and the amplitude of the RF are generated by the computer using a PC30-D card. A simple computer algorithm is used for this purpose. The dependence of the frequency on the VCO input voltage is given in figure A.10.

A.3.1 Digital to analog converter (DAC) summing amplifiers

The interface between the VCO and the computer is formed by the DAC summing amplifier. A detailed schematic is shown in figure A.11. The available 12 bit input-output I/O card gives a resolution of 2.44 mV, which translates to 12 kHz. This does not fulfill the resolution requirements of 6kHz, so we have effectively converted the existing card to a 13 bit card by adding an extra bit. This additional bit was generated using a TTL logic bit from the same I/O card (connector J1 pin 6) and a network of resistors R8-R11. In this way we have increased the I/O card resolution, or equivalently decreased the step size of the output voltage from 2.44 mV to 1.22 mV, which corresponds to a VCO frequency shift of 6 kHz. The integrated circuit U2 is a summing amplifier which produces an
A.3. AOM DRIVERS

output signal for the control of the VCO. The range of the output voltages of the summing amplifiers can be adjusted. This way, the frequency band of the VCO can be adjusted, via potentiometers RV1 and RV3 for the minimum and maximum respectively (in our case 100-150 MHz). The voltage increment set by the TTL bit can be adjusted with potentiometer RV4.

A.3.2 VCO frequency comparator

The VCO frequency comparator allows the measurement of the difference in frequency between two VCOs and is shown in detail in figure A.12. A double balanced passive mixer, SBL1, gives the sum or difference of frequencies of the two VCOs. This is followed by a low pass filter with 10.7 MHz corner frequency, that provides an input signal (the difference of two VCO frequencies) to a comparator U1. Positive feedback to the comparator formed by resistors R4/R3, generates a small hysteresis, minimising the effects of noise on the output signal.

A.3.3 VCO-RF board

The heart of the AOM driver is the VCO-RF board. It consists of a VCO, an analog attenuator, an RF switch and a fixed gain power amplifier. The control signals are derived from the summing amplifier board. A detailed schematic is shown in figure A.13. The signal from the VCO (U1) is AC coupled into an -8 dB attenuator formed by resistors R1-R3, which together with the +8dB amplifier (U2) forms a buffering stage for the VCO. This buffer stage is followed by a power splitter, whose role is to take a sample of the RF to the RF comparator stage. The RF signal is then attenuated by an analog attenuator and can be switched on sub-microsecond timescales by an RF switch. We have decided to use an analog attenuator which allows continuous change of the AOM's amplitude, but there is a provision on the board for an implementation of a digital attenuator, if desired. The attenuator driving circuit, consisting of U9 and Q1, is designed to fulfill current requirements for the attenuator input. U7 is an RF amplifier with maximum output of 2 W.
Figure A.11: Digital to analog converter (DAC) summing amplifiers
Figure A.12: VCO frequency comparator
Figure A.13: The VCO RF Board
References


REFERENCES


[33] T. Bergeman. to be published.

REFERENCES


[56] Generally, di-electric mirrors introduce a substantial phase shift between the s and p polarised components of an incident laser beam. The polarisation preserving mirrors used here were manufactured by Laser Optik Gmbh, Garbsen, Germany, and were specified to introduce no such phase shift. This was verified to be correct. However, the Eindhoven group experienced a substantial phase shift (~ 10 degrees) when the mirrors are placed in vacuum. In their setup no built in compensation was used, and hence the mirrors were virtually unusable. In our setup we have not noticed any adverse effects of the phase error.


REFERENCES


[70] Skytek Corporation. Castech, price quotation 97.

[71] Quantum technology. Private communication.


REFERENCES


