

Title of Dissertation: GRATING MAGNETO-OPTICAL TRAPS
FOR STRONTIUM

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This thesis describes the construction of a new laboratory and a new apparatus for ultracold strontium experiments. It also describes the the creation of the first strontium grating magneto-optical trap (MOT). The new strontium apparatus has been improved to provide high, and controllable atomic flux, high magnetic field coils, low vacuum pressure, plenty of optical access, and a streamlined computer control system. This flexible apparatus is designed to operate several potential experiments.

Our work with strontium grating MOTs will help to enable the creation of compact, stable, and portable quantum devices that harness the unique capabilities of alkaline earth atoms. We present the construction of the grating MOT apparatus, the demonstration of the broad line strontium grating MOT, and the transfer to the narrow line grating MOT. In addition, we show the effectiveness of sawtooth wave adiabatic passage (SWAP) in the grating MOT geometry.

Lastly, we present work from the previous apparatus, measuring the isotope shifts on the clock and intercombination transitions in strontium.

GRATING MAGNETO-OPTICAL TRAPS FOR STRONTIUM

by

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List of Abbreviations

AOM	Acousto-Optic Modulator
AR	Antireflection
CAVS	Cold Atom Vacuum Standards
CG	Center of Gravity
DDS	Direct Digital Synthesizer
DMD	Digital Mirror Display
ECDL	Extended Cavity Diode Laser
EOM	Electro-Optic Modulator
FSR	Free Spectral Range
IGBT	Insulated-Gate Bipolar Transistor
JQI	Joint Quantum Institute
MOSFET	Metal–Oxide–Semiconductor Field-Effect Transistor
MOT	Magneto-Optical Trap
NEG	Non-Evaporable Getter
NI	National Instruments
NIST	National Institute of Standards and Technology
OD	Optical Depth
PBS	Polarizing Beamsplitter
PDH	Pound-Drever-Hall
PID	Proportional Integral Differential
PLL	Phase Lock Loop
RF	Radio Frequency
SRS	Stanford Research Systems
SWAP	Sawtooth Wave Adiabatic Passage
TA	Tapered Amplifier
ULE	Ultra Low Expansion glass
VCO	Voltage Controlled Oscillator

Chapter 1: Introduction

1.1 Strontium Basics

Laser cooling and trapping was initially developed with alkali atoms, however, strontium and other alkaline earth atoms have unique properties that make them useful for a variety of applications, from quantum simulation to precision metrology. Strontium has four stable isotopes, three bosons (^{84}Sr , ^{86}Sr , and ^{88}Sr), and one fermion (^{87}Sr). The properties of these isotopes are summarized in Tab. 1.1. The nuclear spin I is zero for all the bosonic isotopes of strontium and other alkaline earth elements, and thus they lack any hyperfine structure.

The most fundamental change when considering alkaline earth atoms over alkali atoms is the inclusion of a second valence electron. These two electrons can either be in a singlet or triplet configuration, creating more complicated level structures. The level structure of strontium applicable to our experiments is shown in Fig.

	^{84}Sr	^{86}Sr	^{87}Sr	^{88}Sr
Abundance	0.56%	9.86%	7.00%	82.58%
Nuclear Spin	0	0	9/2	0
Scattering Length	$-2a_0$	$96a_0$	$823a_0$	$123a_0$

Table 1.1: Properties of the different stable isotopes of strontium. $a_0 = 5.29 \times 10^{-11}$ m is the Bohr radius.

1.1. To first order, electric dipole transitions don't couple singlet and triplet states, but the spin-orbit interaction mixes states with the same total angular momentum J , allowing these transitions to happen [1]. For instance in strontium the “inter-combination” transition $^1S_0 \rightarrow ^3P_1$ is possible due to mixing between 3P_1 and 1P_1 [2]. This mixing effect also gives these transitions narrow linewidths, making them useful for creating cold magneto-optical traps (MOTs) around a few μK without the need for sub-Doppler cooling.

The spin-orbit interaction does not allow for the transition $^1S_0 \rightarrow ^1P_0$, dubbed the “clock” transition, since $J = 0 \rightarrow J' = 0$ is forbidden for dipole transitions. However, The hyperfine interaction of the fermion, mixes states with the same F , so that 3P_0 state mixes with 3P_1 , 1P_1 , and 3P_2 to allow for the transition [2]. Since this effect is much weaker than the spin-orbit interaction, the linewidth of this transition is very narrow ~ 1 mHz. While the clock transition is generally not allowed in the bosonic isotopes, as they do not have hyperfine structure, applying a magnetic field mixes 3P_0 with 3P_1 , and allows for the transition [9].

In another contrast to alkalis, strontium and other alkaline earth atoms have no electronic angular momentum ($J = 0$) in their ground states. Thus the ground state is insensitive to magnetic fields which eliminates, or complicates, some useful tools such as magnetic traps, Feshbach resonances [10], and sub-Doppler cooling. While ^{87}Sr has $F = 9/2$ in the ground state, it still has a very small g -factor due to the small nuclear magneton. In addition, since the excited state of the clock transition (3P_0) also has $J = 0$, the transition is also magnetically insensitive. Even the small g -factors of these states cancel when considering the differential magnetic

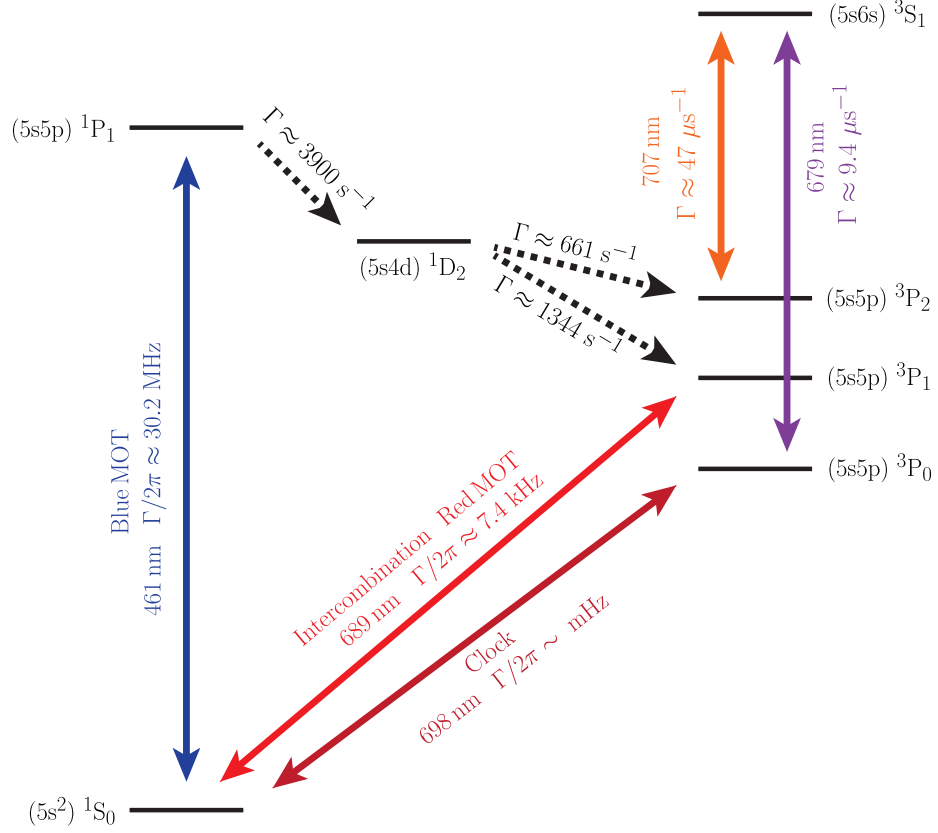


Figure 1.1: Level diagram for strontium. The blue MOT, red MOT / intercombination, and clock transitions from the ground state are shown with wavelengths and linewidths [3–5]. The decay out of 1P_1 state [6], and subsequent decays into the 3P states [7] (from context and alignment with other theses I assumed the 1D_2 - 1P_1 rate listed in [7], was in fact the 1D_2 - 3P_1 rate). The two repump transitions are also shown with their wavelengths and linewidths [8].

sensitivity, and only the contribution coming from the state mixing remains [8]. The narrow linewidth and magnetic insensitivity of the clock transition justify its name as is used for some of the world’s best optical atomic clocks [11–13]. The clock transition also makes strontium useful for different types of metrology and quantum simulation discussed throughout this thesis.

While the bosonic isotopes of strontium have no nuclear spin, the fermion has a relatively high nuclear spin ($I = 9/2$). In addition, the $J = 0$ nature of the ground and clock states decouple the electronic angular momentum from the nuclear spin. This decoupling also makes collisions and atomic interactions independent of the sublevel m_I , which creates an $SU(N)$ spin symmetry in the nuclear spin sublevels, where N can be chosen as long as it satisfies $N \leq 2I + 1 = 10$ for strontium [14, 15]. This large symmetry group makes strontium an attractive element for quantum simulations of systems with this symmetry [16–22].

For reference, I have included a few tables with different properties of strontium. Tab. 1.2 lists the wavelength, linewidth, saturation intensity, and fine structure g -factors of the blue MOT, red MOT / intercombination, and clock transitions. Tab. 1.3 lists the different isotope shifts relative to ^{88}Sr for the main transitions. Tab. 1.4 lists the g -factors for the different excited hyperfine states of ^{87}Sr .

1.2 Isotope Shift Introduction

The final experiment for the old strontium laboratory involved precision measurements of isotope shifts on the intercombination and clock transitions in stron-

	$^1S_0 \rightarrow ^1P_1$	$^1S_0 \rightarrow ^3P_1$	$^1S_0 \rightarrow ^3P_0$
λ	460.9 nm	689.4 nm	698.4 nm
$\Gamma/2\pi$	30.2 MHz	7.4 kHz	~ 1 mHz
I_{sat}	40.4 mW/cm ²	3.0 μ W/cm ²	$\sim 10^{-10}$ mW/cm ²
$g_{J'}$	1	3/2	

Table 1.2: Parameters for different transitions in strontium. The vacuum wavelengths (λ), linewidths ($\Gamma/2\pi$), saturation intensities (I_{sat}), and excited states fine structure g -factors ($g_{J'}$) are listed. Numbers from [3, 5, 8].

Isotope Shift (kHz)	$^1S_0 \rightarrow ^1P_1$	$^1S_0 \rightarrow ^3P_1$	$^1S_0 \rightarrow ^3P_0$
84-88	-273	-351.5	-349.7
86-88	-126	-163.8	-162.9
87-88 (CG)	-45	-62.2	-62.2
87-88 ($F' = 7/2$)	-8	1351.9	
87-88 ($F' = 9/2$)	-68	221.7	
87-88 ($F' = 11/2$)	-50	-1241.5	

Table 1.3: Isotope shifts for the different transitions in strontium relative to ^{88}Sr . For states with hyperfine structure the center of gravity (CG) shift, weighted by number of sublevels, is shown along with individual hyperfine transitions. Isotope shifts are from Refs. [23, 24].

g_F	1P_1	3P_1
$F = 7/2$	-2/9	-1/3
$F = 9/2$	4/99	2/33
$F = 11/2$	2/11	3/11

Table 1.4: Hyperfine g -factors for the excited states of ^{87}Sr . This does not take into account the nuclear magneton.

tium. Precision spectroscopy is a useful tool in general, as it provides the basis for more precise experiments in the future. It can also provide insight into atomic structure, and be used to search for new physics via comparison with theoretical calculations. Isotope shifts in particular give insight into nuclear and atomic structure. In addition, by comparing the isotope shifts on two transitions we can extract a linear relationship, known as King linearity [25].

The King linearity is a useful tool to track the change in the charge radius as neutrons are added to the nucleus and the charge distribution changes. In addition, this linear relationship has been proposed as a way to search for new physics, by putting limits on new physics that would create deviations from this linear relationship [26, 27]. However, higher order effects within the standard model can also cause nonlinearities, so the linearity of King plots only becomes a useful tool when these effects are well accounted for.

Our lab measured the isotope shifts for both the intercombination and clock transitions between four stable isotopes of strontium at the level of ~ 10 kHz. With the four isotopes of strontium we had just enough measurements to investigate King linearity in this system, and found a significant deviation from linearity. This deviation is most likely due to higher order effects in the calculation of the center of gravity of the intercombination line of ^{87}Sr due to its hyperfine structure. Nevertheless, this work provided the first precision measurement of the clock transition in ^{86}Sr and ^{84}Sr , and from our King plot analysis we were able to extract useful information about the mass and field shift parameters for these transitions.

1.3 New Strontium Laboratory

Our work on isotope shifts was the last experiment performed in our old strontium laboratory. While this project was being completed we were simultaneously building a new laboratory space and apparatus and implementing a number of improvements. The new strontium laboratory aims to provide the tools required for a variety of experiments that utilize the unique properties of strontium. We have investigated using Rydberg dressing to increase the precision of optical atomic clocks [28, 29], generating sub-wavelength attractive potentials [30], and using a box trap for quantum simulation [16, 17]. These experiments require a robust, stable, and flexible apparatus.

During my time in grad school we have built up a new laboratory and apparatus with the capabilities to execute these experiments. We implemented a new, high flux, strontium source that has increased the atom number in our MOTs. We designed and built new, high current magnetic coils for operating MOTs, or for other high field investigations. We rebuilt some laser systems, and refined others. Finally, we implemented laser cooling and trapping of strontium, from the 2D MOTs of our source, to the broad blue MOT, to the narrow line red MOT, to dipole traps, which prepare our atomic gases for future experiments.

1.4 Grating MOT Introduction

Many of the unique properties of strontium and other alkaline earth atoms make them particularly useful for metrology and sensors. However many of these applications only reach their full potential with stable, compact, and portable laser-cooled quantum devices. For instance, portable quantum devices with strontium could be used for gravimetry [31] or inertial navigation [32], devices capable of being installed in satellites could be used as gravitational wave detectors [33–35], and accurate dissemination of the SI second relies on the availability of simple and stable atomic clocks [36]. Generally, experiments with laser-cooled atoms, and in particular conventional 6-beam MOTs, require large laboratory spaces, large vacuum chambers, and many laser beams, which makes them difficult to engineer into a stable usable device.

In our laboratory, we worked on developing grating MOTs for strontium as one solution to this issue. By using a single laser beam input and the diffracted beams off of a grating to create the trapping forces, grating MOTs offer simplification, stability, and miniaturization over traditional 6-beam MOTs. Only alkali atoms had previously been trapped in a grating MOT before our work with strontium. We created the first alkaline earth grating MOT with our strontium blue grating MOT, implemented transfer to the narrow line red MOT, and showed sawtooth wave adiabatic passage (SWAP) working in our narrow line grating MOT.

1.5 Outline of Thesis

The following chapters will go into detail on the work I have done over my time in graduate school. Chapter 2 will give an overview of the theory behind the laser cooling and trapping techniques for strontium, grating MOTs, and SWAP. Chapter 3 will give an overview of the new experimental apparatus, both the main apparatus and the grating MOT apparatus. Chapter 4 will discuss the various cooling stages performed on the main apparatus, and their optimizations and results, it will also give a description of the imaging process and the results and process of running the grating MOT. Chapter 5 includes two publications. One publication on our measurements of Isotope shifts, which I supported at the start of the experiment, and took over the for the second half of data collection. The other publication on the blue grating MOT, which was a joint project with me and Ananya Sitaram, and with the guidance of Daniel Barker. Chapter 6 will conclude the thesis with the addition of forward-looking experiments that could be performed on our two systems.

Chapter 2: Theory

2.1 Laser Cooling of Strontium

The experiments and applications that our laboratory aims to accomplish require us to first create a cold sample of atoms. In this chapter I will first go through a relatively simple explanation of the laser cooling techniques used to generate our cold atom gases. Much of the theory for laser cooling is presented with more detail in Ref. [37]. I will then go through each cooling stage, and how they apply to the specific case of strontium in our system.

2.1.1 Basics of laser cooling and trapping

In order to cool and trap atoms with light, we must find a way to apply a force to the atoms. The first force we consider is the scattering force caused from the absorption and emission of photons by atoms. First let us break down the interaction of atoms with photons into three processes: absorption, stimulated emission, and spontaneous emission. A photon carries a momentum of $\vec{p} = \hbar\vec{k}$, where $\vec{k}$ is the wavevector, $k = 2\pi/\lambda$. Thus absorption of a photon from a laser beam transfers momentum to the atom in the beam's propagation direction. Stimulated emission,

the time reversal of absorption, of a photon into the laser beam kicks the atom with the same momentum but in the opposite direction of the laser beam. Thus these two processes together, from the same laser beam, give no net force on the atom, however, the inclusion of spontaneous emission changes this picture. When an atom spontaneously emits a photon the emission is in a random direction, following a dipole distribution, this process provides no net momentum when averaged over many emissions. So when we simply combine these three processes from one laser beam, the only way to get net momentum transfer is absorption followed by spontaneous emission. The rate for this process, for a two level atom, is simply the rate of spontaneous emission: $\Gamma\rho_{ee}$, where ρ_{ee} is the population in the excited state from the density matrix, and Γ is the decay rate. This gives a force of:

$$\vec{F} = \hbar\vec{k}\Gamma\rho_{ee} = \frac{\hbar\vec{k}s_0\Gamma/2}{1 + s_0 + (2\delta/\Gamma)^2}, \quad (2.1)$$

where $s_0 = 2\Omega^2/\Gamma^2$ is the on-resonance saturation parameter, Ω is the Rabi frequency and δ is the (angular) detuning. We now have a simple way to provide a force on atoms in the propagation direction of a resonant laser beam. However, while a constant force on an ensemble of atoms may be able to push it around, in order to cool or trap the ensemble we require a velocity or spatially dependent force respectively.

There are three main sources of line shifts that allow us to create these inhomogeneous forces. Firstly the Doppler shift: here the frequency of light, in the

atom's frame, shifts depending on the atomic velocity as

$$\Delta\omega = -\vec{k} \cdot \vec{v},$$

where $\vec{v}$ is the velocity of the atoms. With Doppler shifts alone you can cool atoms in an optical molasses consisting of counter-propagating red detuned laser beams. Secondly the Zeeman shift: here the energy of an atomic state is dependent on a magnetic field, which for weak fields goes like

$$\Delta E = g_F \mu_B m_F B,$$

where g_F is the g-factor, μ_B is the Bohr magneton, B is the magnetic field, and m_F is the projection of the angular momentum along the magnetic field direction. With the Zeeman shift alone you can create a magnetic trap where a spacial magnetic field gradient creates a trapping potential for atoms in magnetic field sensitive states, discussed briefly in Sec. 2.1.2. With the Zeeman and Doppler shifts together you can create a Zeeman slower [38], where a spatially varying magnetic field is engineered such that, as atoms are slowed with a laser beam, the Doppler and Zeeman shifts cancel to maintain a constant resonance frequency, and thus a constant slowing force. Additionally these two effects together form the foundation of a magneto-optical trap (MOT) which will be discussed in Sec. 2.1.2, and 2.1.3. The last shifts to consider are light shifts: here the presence of an oscillating electric field far detuned from a resonance gives a shift, in a two level system, to the ground (g) and

excited (e) states of

$$\Delta E_{g,e} = \pm \frac{\hbar \Omega^2}{4\delta} = \pm \frac{\hbar \Gamma^2 I}{6 I_s \delta}, \quad (2.2)$$

where Ω is the Rabi frequency, proportional to the magnitude of the electric field, I is the intensity of the light, and I_s is the saturation intensity. Thus for a laser beam detuned below resonance, atoms in the ground state will be attracted to locations of high intensity. In this way tightly focused laser beams can create a trap for atoms called an optical dipole trap discussed in Sec. 2.1.4. All of these cooling and trapping techniques depend on the specifics of the atomic element used. In the following sections I will discuss how these techniques apply to strontium specifically.

2.1.2 Blue MOT

A standard six-beam MOT consists of a quadrupole magnetic field and six counter-propagating, appropriately circularly polarized, red detuned, orthogonal, laser beams. Our “blue” MOT of strontium operates on the $^1S_0 \rightarrow ^1P_1$ transition (see Fig. 1.1). We will first consider the bosonic isotopes where there is no hyperfine structure, which simplifies the picture. Also, since the ground state has $J = 0$ we only have to consider the Zeeman shift to the excited state sublevels. To consider the force profile in the MOT we can use the scattering force from Eq. 2.1, and include the relevant shifts to the transition in the detuning. If δ is the detuning from the bare resonance, then including the Doppler and Zeeman shifts, the detuning from

the shifted resonance becomes:

$$\delta' = \delta - \vec{k} \cdot \vec{v} - \mu_B m_F B / \hbar,$$

where $g_F = 1$ for the 1P_1 state. We consider the scattering force in only the x dimension of the MOT where the magnetic field has a linear gradient of $\vec{B} = dB_x \hat{x}$. Here we combine the force from the two counter-propagating beams, one with $\vec{k} = +\hat{x}$ addressing the $m_F = +1$ excited state, and one with $\vec{k} = -\hat{x}$ addressing the $m_F = -1$ state to get

$$F_x = \hbar k s_0 \frac{\Gamma}{2} \left(\frac{1}{1 + s_0 + 4(\delta - kv_x - \mu_B dB_x x)^2 / \Gamma^2} - \frac{1}{1 + s_0 + 4(\delta + kv_x + \mu_B dB_x x)^2 / \Gamma^2} \right).$$

Thus with red-detuned beams we get a spatial and velocity dependent force that slows atoms, and traps them near $x = 0$. The force in equation 2.1 is in reality an average over many photon kicks, so while an atom at the center of the MOT with zero velocity may see no average force, there is still a non-zero variance to the force, which causes heating in the MOT. The Doppler temperature is the lowest temperature possible with Doppler theory due to this heating, and is defined as

$$T_D = \frac{\hbar \Gamma}{2k_B},$$

where k_B is the Boltzmann constant, and $\Gamma/2\pi = 30.2$ MHz. For the blue MOT this temperature is 840 μ K. While there are many different complications that increase

the MOT temperature, or techniques to get below this limit, the Doppler limit gives us a rough idea of how cold atoms in our blue MOT will be.

While the blue MOT is simple compared to the following red MOT cooling stage, there is one extra complication coming from the level structure. The excited state of the cycling transition, 1P_1 , can decay through the 1D_2 state into the 3P_1 and 3P_2 states as shown in Fig. 1.1. The 3P_1 state decays relatively quickly, with a lifetime around $20\ \mu\text{s}$, but the 3P_2 state does not decay, requiring repump lasers to excite atoms out of this state and keep them cycling, and therefore cooling. There are many different repumping schemes for strontium [39–42], but the one we use requires two lasers. The first at 707 nm pumps atoms out of the 3P_2 state to the 3S_1 state. From there they can decay to the 3P_1 state or to the metastable 3P_0 state. This second state is the excited state of the clock transition, with a very long lifetime, during which atoms will be lost from the trap. To avoid this we use our second repump laser at 679 nm to pump out of this state to the same 3S_1 state. Thus any atoms in the blue MOT that become dark to the cycling transition are pumped through 3P_1 back to the ground state.

While the metastable states complicate the blue MOT slightly, they also open up an opportunity for magnetic trapping. While the ground 1S_0 state cannot be magnetically trapped due to the lack of angular momentum sublevels subject to a Zeeman shift, the 3P_2 state can be magnetically trapped. Even without repumps, some of the atoms that decay to the 3P_2 state land in sublevels that are magnetically trapped by the quadrupole field of the MOT. This leads to techniques where, while loading the blue MOT, you can leave the repump lasers off to load this magnetic

trap of $^3\text{P}_2$ atoms, called the metastable reservoir, and pulse on the repumps at the end of the blue MOT cooling stage to reintroduce these atoms to the system, hopefully increasing the number of available atoms in total. We don't use this technique often, but it can be useful for the less abundant isotopes. Due to the long lifetimes of the metastable states, running the magnetic trap allows for increased accumulation these less abundant isotopes compared to the blue MOT. More details can be found in Ref. [43].

So far we have been ignoring any hyperfine structure, which is appropriate for the bosonic isotopes, but not for the fermion: ^{87}Sr . Luckily, the blue MOT doesn't require much alteration, other than accounting for isotope shifts. The only other change is to account for the additional hyperfine structure in the 707 nm repump transition by broadening that laser in an attempt to address all of the new hyperfine levels. More details on this broadening are described in Sec. 4.2.

The blue MOT allows us to cool and trap atoms out of our source, but the cooling is not sufficient to load into a dipole trap or work with the clock transition, so instead, we transfer to a second, colder MOT.

2.1.3 Red MOT

The next stage of cooling is a MOT operating on the narrow $^1\text{S}_0 \rightarrow ^3\text{P}_1$ “inter-combination” line. The linewidth of this transition is just 7.5 kHz, meaning that the Doppler limit is about 1.8×10^{-7} K. While the narrow linewidth of the red transition helps to achieve a colder temperature, it also has the side effect that light at a single

frequency is only resonant with a small shell of atoms in phase space. In addition, the lower forces in the red MOT compared to the blue MOT cause gravity to become significant, and the atoms fall to sit at the bottom of this resonant shell. Since the blue MOT is relatively hot, it has a wide velocity distribution, and expands rapidly. Thus the resonant shell of atoms that the red light can address is a very small percentage of the total blue MOT population. In order to achieve a sufficient transfer efficiency between the MOTs we initially broaden the red cooling laser to address as many velocity classes and positions as possible. This broadening is ramped down over time to compress the MOT and cool to our final temperature. A more specific description of this transfer with our experimental parameters is presented in Sec. [4.3](#).

Where the red MOT really becomes complex is with the addition of the hyperfine structure in ^{87}Sr . Given the nuclear spin $I = 9/2$ our red MOT operates on a $F = 9/2 \rightarrow F' = 11/2$ transition. There are two properties of strontium that complicate the fermionic red MOT: first is the negligible magnetic sensitivity of the $J = 0$ ground state, and second is the narrow linewidth which fully resolves the hyperfine structure. In the simple MOT description from the previous sections the shift of σ^+ transitions as a function of field have positive slopes, and vice versa for σ^- . This allows for selective excitation from light with the appropriate k -vector. As long as g_F and $g_{F'}$ are close enough in value, adding in hyperfine structure will not change this picture much, as shown for ^{87}Rb in Fig. [2.1](#) (b). In particular, as long

as the σ^+ transitions for all m_F shift such that

$$g_{F'}\mu_B(m_F + 1) - g_F\mu_B m_F > 0,$$

the MOT can operate stably. Considering the stretched states $m_F = \pm F$ this requirement becomes:

$$\frac{F}{F+1} < \frac{g_{F'}}{g_F} < \frac{F}{F-1}$$

However in the ^{87}Sr red MOT transition, shown in Fig. 2.1 (c),(d), due to the lack of magnetic sensitivity in the ground state, these conditions are not met. The shift in σ^+ transitions from $m_F < 0$ states have a negative slope with magnetic field. Thus, depending on the m_F state and position of an atom, it can be resonant with the wrong laser beam and be kicked out of the trap. In addition for $m_F \geq 3/2$ ($m_F \leq -3/2$), there is no resonant light at all on the $+x(-x)$ side of the trap.

If we can keep atoms in the $m_F = \mp F$ states on the side of the trap with $\pm$ magnetic fields we can avoid these anti-trapping effects in two ways. First, an atom in one of these appropriate states will be resonant with the correct laser beam at a smaller displacement than the resonant position of the anti-trapping beam. Second, the Clebsch-Gordon coefficients from these appropriate sublevels favor the trapping transitions over the anti-trapping transitions as shown in Fig. 2.1 (d). Not only does this mean the trapping transitions are stronger, but it also means that atoms are pumped into the ideal stretched $m_F = \pm 9/2$ state, where the trapping transition is strongest. Unfortunately, this pumping becomes undesirable as soon as the atom

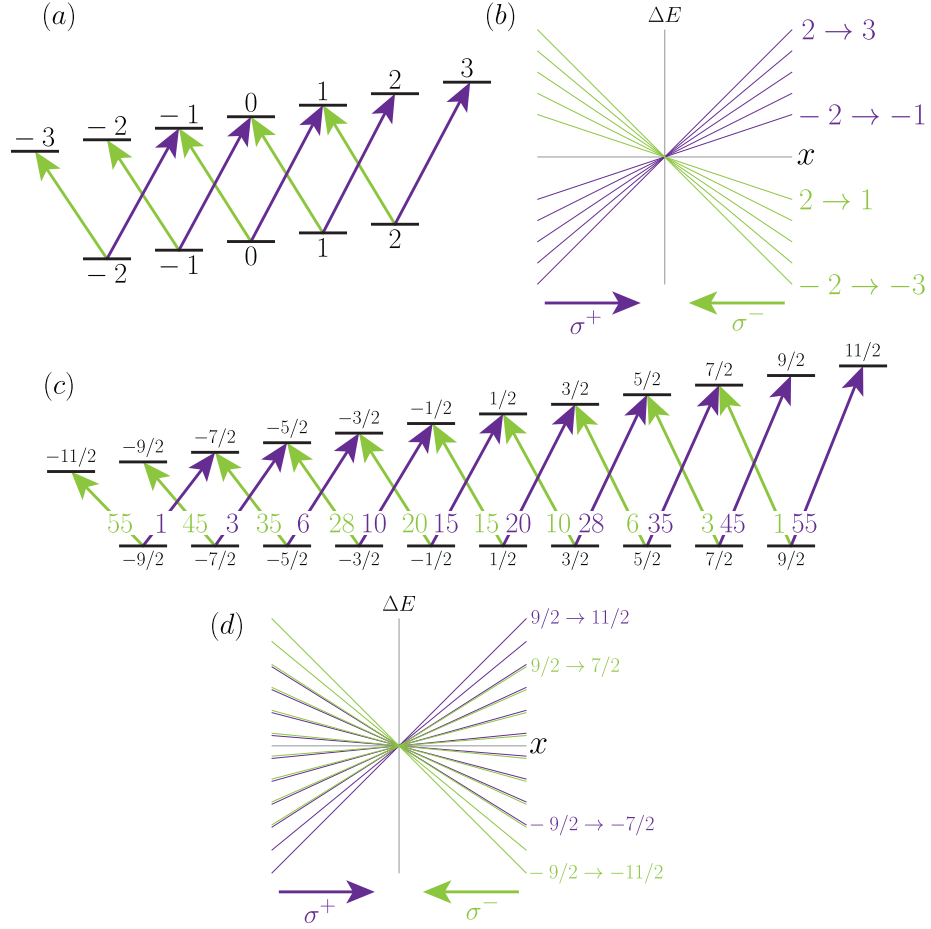


Figure 2.1: Effect of hyperfine structure on the red MOT performance. (a) For comparison, the level structure of the $F = 2 \rightarrow 3$ transition in ^{87}Rb , and its level shifts from a positive magnetic field. (b) The level splittings from (a) in a field gradient. For negative detuning you can see the σ^+ transitions are grouped together all in the same quadrant, and isolated from the σ^- transitions, and vice versa. (c) The level structure of the red MOT transition $F = 9/2 \rightarrow 11/2$ in ^{87}Sr , and the shifts from a positive magnetic field. The Clebsch-Gordan coefficients, normalized to integer values, are shown on their respective transitions. For strontium, the ground state field shift is negligible. (d) The level splitting from (c) in a field gradient. Some σ^+ transitions have a negative slope, which are anti-trapping, and vice versa for σ^- . The slight offset between certain σ^+ and σ^- transitions are simply for clarity and are not to scale.

crosses to the other side of the trap, where it will be in the wrong stretched state and will not see any resonant laser beams at all.

While the arguments made above apply to the fermionic blue MOT as well, that MOT works well enough because of the much larger linewidth. In the blue MOT, the hyperfine structure is not resolved, so the MOT beams addressing the $F = 9/2 \rightarrow 11/2$ cooling transition, are also addressing the $F = 9/2 \rightarrow 9/2$ transition. This added excitation randomizes the populations in the ground state hyperfine manifold, negating the negative effect of the optical pumping, we call this effect “stirring”. For the red MOT transition the linewidth is much smaller than the hyperfine splittings, so the hyperfine structure is fully resolved, and we do not see the stirring effect from the cooling light. However, we can simply add in a stirring laser resonant with the $F = 9/2 \rightarrow 9/2$ transition to replicate this effect from the blue MOT and stabilize the trapping in the fermionic red MOT [44]. Details about the implementation of this are discussed in Sec. 4.3.

2.1.4 Dipole Traps

The atoms in the red MOT are cold enough ($2 - 5 \mu\text{K}$) to be captured directly in a dipole trap. To analyze the potential from this trap we need to understand the intensity profile of the beam used. First consider a single high power laser beam focused to a Gaussian waist of w_0 (typical values range from 10-300 μm), and peak intensity I_0 detuned below resonance ($^1\text{S}_0 \rightarrow ^1\text{P}_1$). Then the transverse intensity profile of the beam is $I(r) = I_0 e^{-2r^2/w_0^2}$. From the light shifts (Eqn. 2.2), this gives

us a transverse potential for atoms in the ground state of

$$V(r) = \frac{\hbar\Gamma^2 I_0 e^{-2r^2/w_0^2}}{6I_s\delta}.$$

Now consider the beam along the axial direction, with a waist of w_0 at $z = 0$, the beam diameter varies like $w(z) = w_0\sqrt{1 + (\frac{\lambda z}{\pi w_0^2})^2}$ [45]. So the intensity, and potential along the axial direction at $r = 0$ become

$$I(z) = \frac{I_0}{1 + (\frac{\lambda z}{\pi w_0^2})^2}, \quad V(z) = \frac{\hbar\Gamma^2 I_0}{6I_s\delta(1 + (\frac{\lambda z}{\pi w_0^2})^2)}.$$

Thus in both the axial and radial directions we have a trapping potential for the ground state with negative detuning, although since the axial trapping is weak often multiple crossed beams are used. In practice it is useful to approximate these potentials as parabolas, so that atoms in the trap can be modeled as simple harmonic oscillators, and we can write the potential as

$$V = \frac{1}{2}m\omega_x^2 x^2 + \frac{1}{2}m\omega_y^2 y^2 + \frac{1}{2}m\omega_z^2 z^2.$$

This treatment is appropriate for a two level atom, but in reality we have to consider contributions to the light shift from many different transitions.

Dipole traps allow for the tight confinement and cold temperatures necessary to perform many experiments, but one issue arises once you try to address transitions in the trap, especially for the narrow transitions in strontium. The potential described above depends on the state of the atom, specifically the AC polarizability, which is

a function of wavelength. So atoms in an excited state will see a different potential than ground state atoms, and consequently the transition frequency will vary over the spacial extent of the trap, causing inhomogeneous broadening. In particular, while the ground state, and the $^3\text{P}_{0,1}$ states of strontium are all trapped by a 1064 nm dipole trap, the differential light shift causes the resonance frequency between them to be dependent on the motional state of the atom. A detailed analysis on the effect is described in the appendix of Ref. [24] (Sec. 5.1). In order to avoid this effect, and homogeneously address the narrow transitions of strontium we can instead pick a “magic” wavelength for our dipole trap where the polarizabilities, and thus the trapping potential, of the two states considered are equal. For the clock transition this wavelength is 813 nm [8].

2.2 Isotope Shifts

The energy levels of a given element will shift between different isotopes due to nuclear effects. This effect is important to take into account when laser cooling and trapping multiple isotopes, and can also be used to probe the nuclear structure, and put limits on new physics which motivated our work in [24] (Sec. 5.1).

The isotope shifts can be split into two parts. The first is the “mass shift” which comes from the change in the nuclear mass. The mass shift is further split into two parts, the normal mass shift, from the change in the reduced mass of the electron, and the specific mass shift which arises from electron correlations in multi-electron atoms and is difficult to calculate. Put together, the mass shift between isotope A

and A' is proportional to $1/\mu_{AA'} = 1/m_{A'} - 1/m_A$, where m_A is the nuclear mass of isotope A . The second part of the isotope shift is the “field shift” which arises from the change in the charge radius of the, and thus is proportional to $\delta \langle r^2 \rangle_{AA'}$ [25], the change in the mean square radius of the nuclear charge distribution. We can then write the isotope shift of a transition (i) between two isotopes (A and A') as a combination of these shifts.

$$\delta\nu_i^{AA'} = K_i/\mu_{AA'} + F_i\delta \langle r^2 \rangle_{AA'} , \quad (2.3)$$

where K_i is a constant associated with the mass shift, F_i is the field shift constant, and $\delta\nu_i^{AA'}$ is the change in the transition frequency. Here we have considered the isotope shift of a transition which is simply the sum of the isotope shift of each state involved.

You can extract a linear relationship between isotope shifts of different transitions, which forms the basis for our work in [24] (Sec. 5.1). Consider that the mass and field shift constants are independent of the isotope, and the factors they are associated with ($\mu_{AA'}$ and $\delta \langle r^2 \rangle_{AA'}$) are independent of the transition. This allows us to rewrite Eq. 2.3 in terms of the isotope shift of another transition j

$$\mu_{AA'}\delta\nu_i^{AA'} = \left(K_i - \frac{F_i}{F_j}K_j\right) + \left(\frac{F_i}{F_j}\right)\mu_{AA'}\delta\nu_j^{AA'}$$

This linear comparison between two transitions causes the nuclear charge radius, which is difficult to calculate, to drop out, and allows us to extract information

about the field and mass shift constants. In addition, looking for deviations from this linear relationship has been proposed as a method to search for new physics beyond the standard model [26, 27].

2.3 Grating MOT Theory

While traditional six-beam MOTs are nicely symmetric and simple, they require good optical access and have many degrees of freedom in alignment and polarization. One way to simplify the alignment and make the apparatus more compact is to use a grating MOT instead. Grating MOTs use diffraction gratings to trap atoms between a single input beam, and the diffracted beams off the gratings [46]. In this section I will focus on the style of MOT we have in the laboratory: a tetrahedral grating MOT, that has three gratings (or grating sections) oriented in a plane, and angled $2\pi/3$ relative to each-other, as shown in Fig. 2.2. The diffracted beams and the input beam then combine to create a tetrahedral trapping region. The tetrahedral geometry uses the minimum number of beams, and creates a relatively large trapping region, but other groups use gratings with four sections [48], regardless, the basic ideas of this section will still hold. The most significant difference between a conventional six beam MOT, and a grating MOT relates to the polarizations of the diffracted beams and the magnetic field vector. Since the k -vector of the diffracted beams don't align well with the magnetic field, the polarization of the diffracted beams in the quantization basis of the atoms is no longer circularly polarized in the correct direction, and instead the atom sees all polarizations. This decreases the

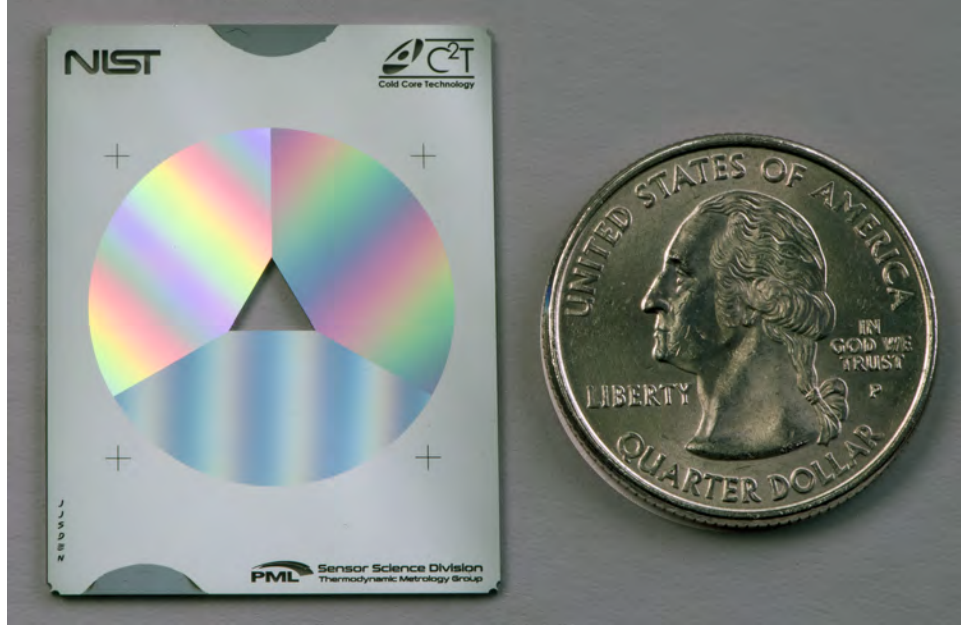


Figure 2.2: A photo of the grating chip used to make our MOTs. Photo from Ref. [47]

effective power of the correct polarization, adds power to the anti-trapping polarization, and adds π polarization as well dependent on the angle of diffraction. For instance, with σ^+ input polarization, where the quantization axis is defined along $\hat{z}$, the fraction of each polarization diffracted from the chip is shown in Fig. 2.3. Here you can see that primarily the anti-trapping polarization, in this case σ^+ is diffracted from the grating. Additionally, the total power in the diffracted beams is significantly lower than the input beam due to the efficiency of the grating. The decreased power, and odd angles of the diffracted beams cause grating MOTs to have anisotropic temperatures [49], and we see this effect when measuring the axial vs transverse temperatures in our MOTs.

We will define our MOT such that the input beam is pointing in the $-\hat{z}$

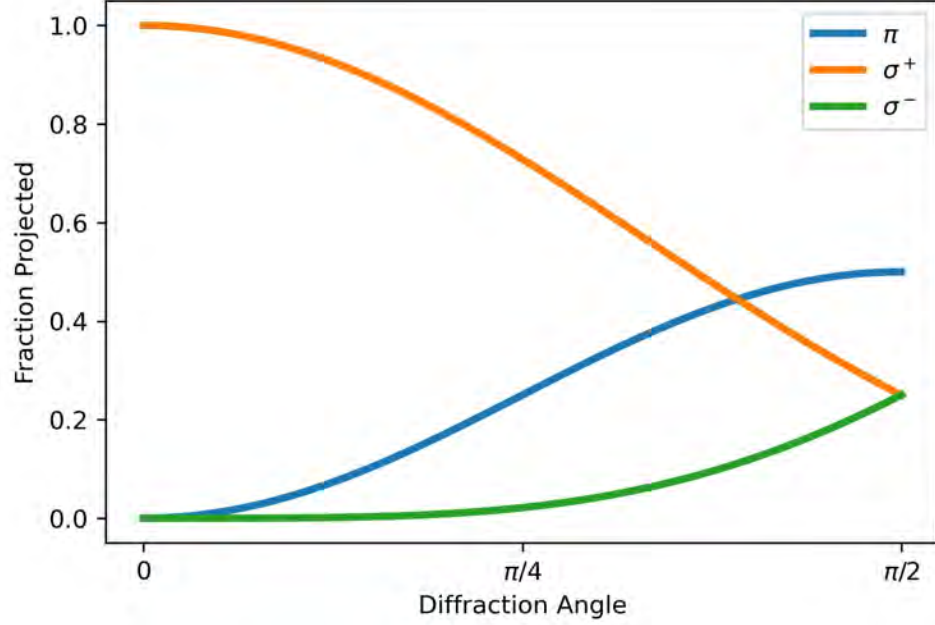


Figure 2.3: The fraction of each diffracted polarization as a function of diffraction angle for a σ^+ input polarization. This is for the quantization axis along $\hat{z}$, where the angle is defined relative to the normal of the grating chip.

axial direction, and the grating is normal to that beam. Taking $\delta = -1\Gamma$, $s_0 = 1$, $\frac{dB_z}{dz} = 55$ G/cm for our blue MOT grating MOT, we can see the force profile in the axial and transverse directions in Fig. 2.4. These force profiles were created with pylcp, a python package for simulating laser cooling physics [50]. The transverse forces in x and y look similar to what we would expect for a traditional MOT, although there is a slight asymmetry caused by the three-fold rotational symmetry of the grating. The axial (z) direction, however looks significantly different from a traditional MOT. The large dip on the $+z$ side of the force profile is due to the force from the input beam, but there isn't a matching peak on the $-z$ side from the diffracted beams. This is due to the lack of power in the correct polarization.

Instead of the traditional circular polarization, our trapping from the diffracted beams comes from the π polarized light which provides force in the the $+z$ direction, but is insensitive to magnetic field, so provides that force everywhere in the beam path. However, the input beam is strong enough to overcome this constant force and provide the force downwards. Consequently, this means that the zero point of the force does not sit at the zero of the magnetic field, and its location depends on a number of atomic, and experimental parameters. In particular, for ^{87}Sr , due to the hyperfine structure and the mismatch in magnetic sensitivity between the ground and excited state, the force from the π polarized light becomes magnetically sensitive. This shifts the force profile so that the fermionic Sr blue MOT forms closer to the field zero than the bosonic MOTs.

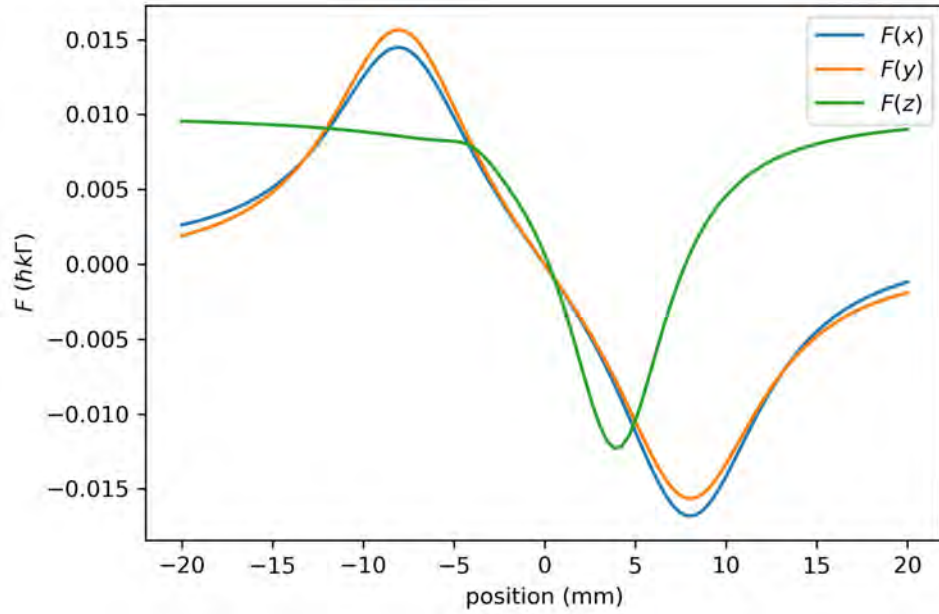


Figure 2.4: The force profiles in the blue grating MOT in x , y , and z , for $\delta = -1\Gamma$, $s_0 = 1$, $\frac{dB_z}{dz} = 55$ G/cm.

Now we can look at the red MOT force profile shown in Fig. 2.5, for $\delta = -10\Gamma$, $s_0 = 90$, $\frac{dB_z}{dz} = 8$ G/cm. Again, it's the axial force that looks significantly different. Here we do see the force from the correct polarization of the diffracted beams, but the peak is much narrower. We see a peak because the saturation intensity of the red MOT transition is much smaller than the blue MOT's, so the power in the diffracted beam is large enough to provide a traditional force. However, the difference in power between the input and diffracted beams is still large, causing power broadening of the input beam's resonance compared to the diffracted beam's resonance.

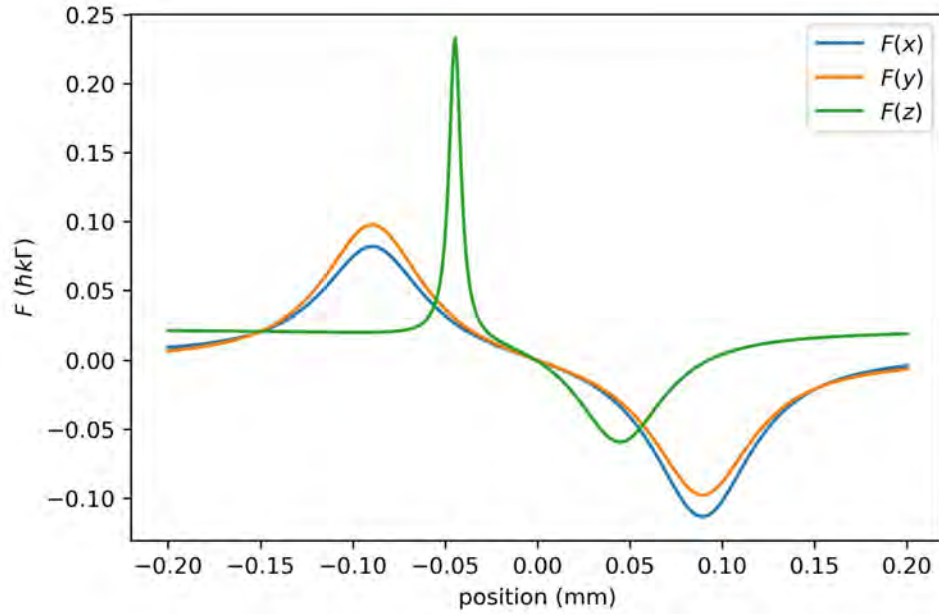


Figure 2.5: The force profiles in the red grating MOT in x , y , and z , for $\delta = -10\Gamma$, $s_0 = 90$, $\frac{dB_z}{dz} = 8$ G/cm.

2.4 Sawtooth Wave Adiabatic Passage

Since the transition between the blue and red MOTs causes significant atom loss, especially in our grating MOT where we lose around 85% of the blue MOT atoms, we have investigated using Sawtooth Wave Adiabatic Passage (SWAP) [51, 52] in the grating MOT to increase the transfer efficiency. SWAP is a cooling technique that works by providing a force against atomic motion from both the absorption and stimulated emission of counter-propagating beams through adiabatic passage. In SWAP two counter-propagating linear polarized beams are swept from below atomic resonance to above it, following a sawtooth function as shown in Fig. 2.6. Consider an atom, starting in the ground state and moving along the axis of the

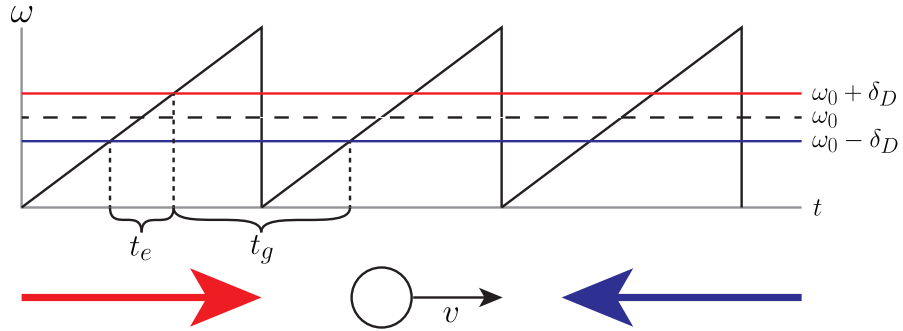


Figure 2.6: The frequency, ω , of the SWAP lasers as a function of time. The bare resonance ω_0 and the frequency shifted by the Doppler shift δ_D are shown. The time that the atom is in the excited state (t_e) and the ground state (t_g) are shown below the time axis. Below the plot is a diagram of an atom with a particular velocity v interacting with the counter-propagating beams.

laser beams. From the Doppler shift the atom first becomes resonant with the beam opposing its motion, which is blue shifted. As this beam sweeps over the resonance

it adiabatically transfers the atom to the excited state, and slows the atom down by $\hbar k/m$. Then as the sawtooth wave continues above the unshifted resonance, the laser co-propagating with the atom is Doppler shifted to the red onto resonance and adiabatically transfers the atom back to the ground state, where the emitted photon again slows the atom down. Then the sawtooth function of the sweep quickly falls back to the start, repeating the process. Compared to traditional Doppler cooling where stimulated emission works against the cooling, the time ordering in SWAP ensures that both absorption and stimulated emission slow the atom.

The atom has to start in the ground state for the time ordering of SWAP to work. If an atom starts in the excited state during a period of SWAP, the reverse process will occur and the atom will be sped up. We can avoid this by making spontaneous emission less likely in between the two transfers, and more likely at the start or the end of the ramp. Thus the time between the transfers (t_e), where the atom is in the excited state, should be less than the lifetime of the state, and less than the total time in the sweep outside of the transfers (t_g). The last requirement [51]

$$\frac{\Omega^2}{\alpha} \geq 1.$$

ensures that the frequency sweep rate α (Hz/s) is slow enough compared to the Rabi frequency of the beams Ω to excite the atoms adiabatically.

Implementing SWAP in a MOT has additional complications. Since the counterpropagating beams in a conventional MOT are σ^+ and σ^- they don't address the same sublevels. Thus, for a simple $F = 0 \rightarrow 1$ transition, atoms adiabatically

excited by one laser beam will not be adiabatically transferred to the ground state by the other laser beam, reducing the efficiency of SWAP. With the right parameters, SWAP can still provide enhancement by suppressing stimulated emission by the wrong laser beam [53]. When operating this way it becomes even more important for the sweep period to be longer than the lifetime since atoms have to decay through spontaneous emission every SWAP period to avoid anti-cooling. Alternatively, one can switch the polarizations of the beams in between the absorption and stimulated emission to maintain the full cooling efficiency of SWAP [54].

The complexities of polarization in a grating MOT ameliorate the issues with performing SWAP in a grating MOT. Grating MOTs have mixed polarizations in all directions (other than the main input beam), so even while operating the MOT there is light at the right polarization to adiabatically transfer atoms to the ground state and impart a cooling force. Fig. 2.7 shows the cooling force in the red grating MOT as a function of atomic velocity for a traditional triangle wave modulation, and sawtooth wave modulation.

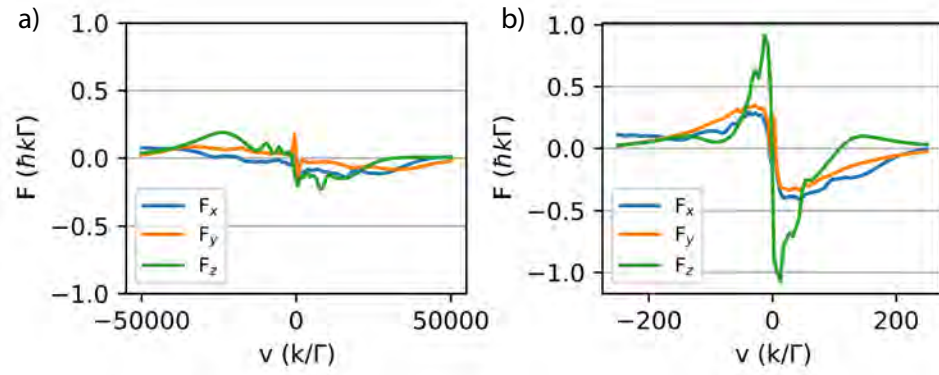


Figure 2.7: Modulated red MOT cooling forces. a) The force vs velocity along all three axes for the red grating MOT with triangle wave modulation. b) The force vs velocity for the MOT with SWAP. You can see the large enhancement in the cooling forces along all directions when performing SWAP.

Chapter 3: Experimental Apparatus

The experimental apparatus involved in laser cooling and trapping experiments are often complicated, interconnected, and expansive, and our laboratory is no exception. Our apparatus is a combination of new and improved elements, and old elements transferred directly from our old laboratory. In this chapter, I will discuss the development of the new vacuum systems, our new magnetic field coils, the variety of laser systems, the optics on the chamber side, and the implementation of computer control throughout the apparatus. I will provide more detail on the systems that we redesigned for the new laboratory; Refs. [5, 55, 56] provide details on the pieces more directly taken from the old apparatus.

3.1 Vacuum System

Neal Pimenti completely redesigned the vacuum system for the new laboratory to avoid the issues present in the design of the old vacuum system, including the large size, lack of optical access, and issues with strontium coatings on windows. I designed and built a high-temperature vacuum bakeout chamber to ensure we were able to attain optimal pressure for this new chamber. In addition, we integrated a high flux commercial source into the new design. The old system is described in

Refs. [5, 55, 56].

3.1.1 Bakeout Chamber

Working with ultra-cold gases requires a vacuum system in the 10^{-10} – 10^{-11} Torr range. Lower pressure increases the lifetime of the trapped atoms and also improves the atom number. However, the rate of desorption of water and hydrogen out of the materials that make up your chamber, in our case polished stainless steel, limits the ultimate pressure of the chamber. Thus it is necessary to bake your vacuum parts at as high a temperature as possible to increase the rate of desorption and pump out the water and hydrogen quickly. Our bakeout process consists of two stages. First, before assembly, the parts are baked to a high temperature (over 400 °C) in a separate vacuum chamber to desorb the hydrogen from the bulk of the material. Second, the whole system is baked at a lower temperature after it is assembled and pumped down to desorb the water at the surface. Both bakes help lower the pressure since the low temperature does not efficiently desorb hydrogen from the bulk of the material, while parts of our vacuum apparatus, such as the viewports, cannot be baked at the high temperature because of the glass to metal seal. The high temperature is particularly important because the rate of desorption approximately doubles for each 10 °C increase in temperature. Additionally, baking in a vacuum will not create an oxide layer on the surface of the stainless steel parts, which would impede the desorption. When the old vacuum chamber was assembled, the laboratory could not perform a high-temperature vacuum bake. Although that chamber had a final

pressure of $\sim 5 \times 10^{-11}$ Torr, we wanted to do a high-temperature vacuum bake for the new chamber to push the limits of the final pressure.

We designed and built a separate vacuum chamber large enough to accommodate each part and capable of reaching high temperatures. Fig. 3.1 shows the bakeout chamber, it is made of a 14 in nipple with a valve to attach a turbo pump.

We use a chicken wire enclosure as a frame for heater tape and insulation. This enclosure allows us to hold the temperature and pressure at around 425 °C, and 1.5×10^{-7} Torr respectively for weeks at a time, monitored by several external thermocouples to control the heating and one vacuum thermocouple to monitor the temperature at the parts that are being baked. We had a proportional integral differential (PID) controller to actively stabilize the temperature of the bake, but due to coupling between the different sections of heater tape, we found manual control to be more effective. This bakeout chamber has become a useful resource throughout the Joint Quantum Institute (JQI) and has since been used by many groups.

The second bake after assembly brought our chamber to ~ 150 ° C and was monitored with a residual gas analyzer and an ion gauge. The pressure of our main chamber is currently around 1×10^{-11} Torr, near the base pressure of our ion gauge, where the reading is not reliable. For example, we discovered during the second bake that the humidity around the ion gauge cable would affect the reading rather significantly.



Figure 3.1: A photo of the full size assembled bakeout chamber with the insulation frame.

3.1.2 AOSense Strontium Source

The strontium source and initial cooling stages on the old chamber had a few issues that we wanted to address in the new laboratory. The old apparatus had a large oven, a transverse cooling stage, and a long Zeeman slower. Together with the main science chamber, the whole vacuum system took up almost an entire optical table. In addition, the Zeeman slower was directly opposite a primary viewport of the science chamber. Over time, strontium coated this viewport, reducing the efficiency of the Zeeman slower and limiting the use of the viewport for other laser beams. To address this issue, the old laboratory heated this viewport to reduce the coating rate, but due to the size of the viewport and proximity to the main science chamber, this solution was not optimal. The most substantial upgrade to the new vacuum system is the commercial source, custom-made by AOSense, shown in Fig. [3.2](#). This source has a much smaller footprint and solves the issue of the strontium coating.

The oven in the source is a small effusive oven, backed with an ion pump. Since the oven is heated in vacuum, it can safely ramp up from room temperature to operating temperature in under an hour, less than half the time of the previous source. There is a short (< 7 cm) permanent magnet Zeeman slower after the oven. This slower is a substantial improvement over the previous slower which was 35 cm long and required additional water cooling. With our new source, the small Zeeman slower window is heated to a constant temperature of 320°C with in-vacuum heaters, and the window is not part of the science chamber itself (see Fig. [3.2](#)).

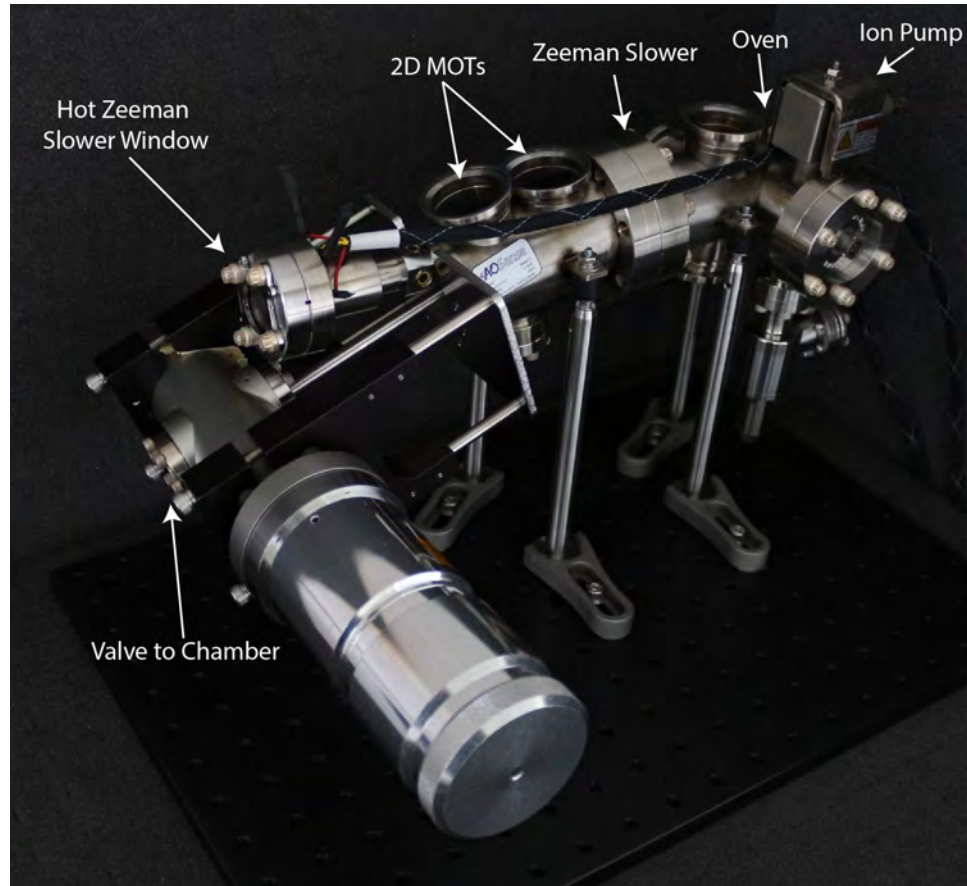


Figure 3.2: A photo of the AOSense source before it was attached to the science chamber. you can see the total size of the source including the minimal length of the Zeeman slower. The progressive angle downwards is also visible, which allows for the heated window the Zeeman slower beam goes through to be separate from the science chamber.

This avoids the strontium viewport coating that occurred in the previous apparatus. Following the slower are two, 2D MOTs that progressively angle the atomic beam downward to 20° from horizontal, and into the vacuum chamber. These 2D MOTs have completely internal optics and magnets, and require only a single input laser beam each. Additionally, the angle means that there will be no flux of atoms into the science chamber unless the 2D MOTs are on, allowing for an effective atomic shutter, replacing our old pneumatic shutter. The tube connecting the source to the main chamber is 10 cm long has an inner diameter of .6 cm, providing differential pumping and maintaining the low pressure of the main chamber. The oven provides an atomic flux of $3.4 \times 10^{12} \text{ s}^{-1}$, at our current operating temperature of 420°C , but can be operated at up to 540°C for a higher flux. Other groups who use AOSense sources have had issues with clogs, they advised us to ramp our source all the way down to room temperature at the end of the day, and that operating at an intermediate temperature will increase the chance of a clog. Since the oven reaches its operating temperature so quickly, operating at an intermediate value is not necessary.

3.1.3 Main Chamber

The final issue with the old chamber was the lack of optical access. One main viewport was coated with strontium, one main port to the chamber was used for the input of the atomic beam, and the mini viewports were blocked by the bucket windows on the top and bottom, rendering them useless. Our new science chamber,

shown in Fig. 3.3, is similar to our previous chamber, but with significantly more optical access.

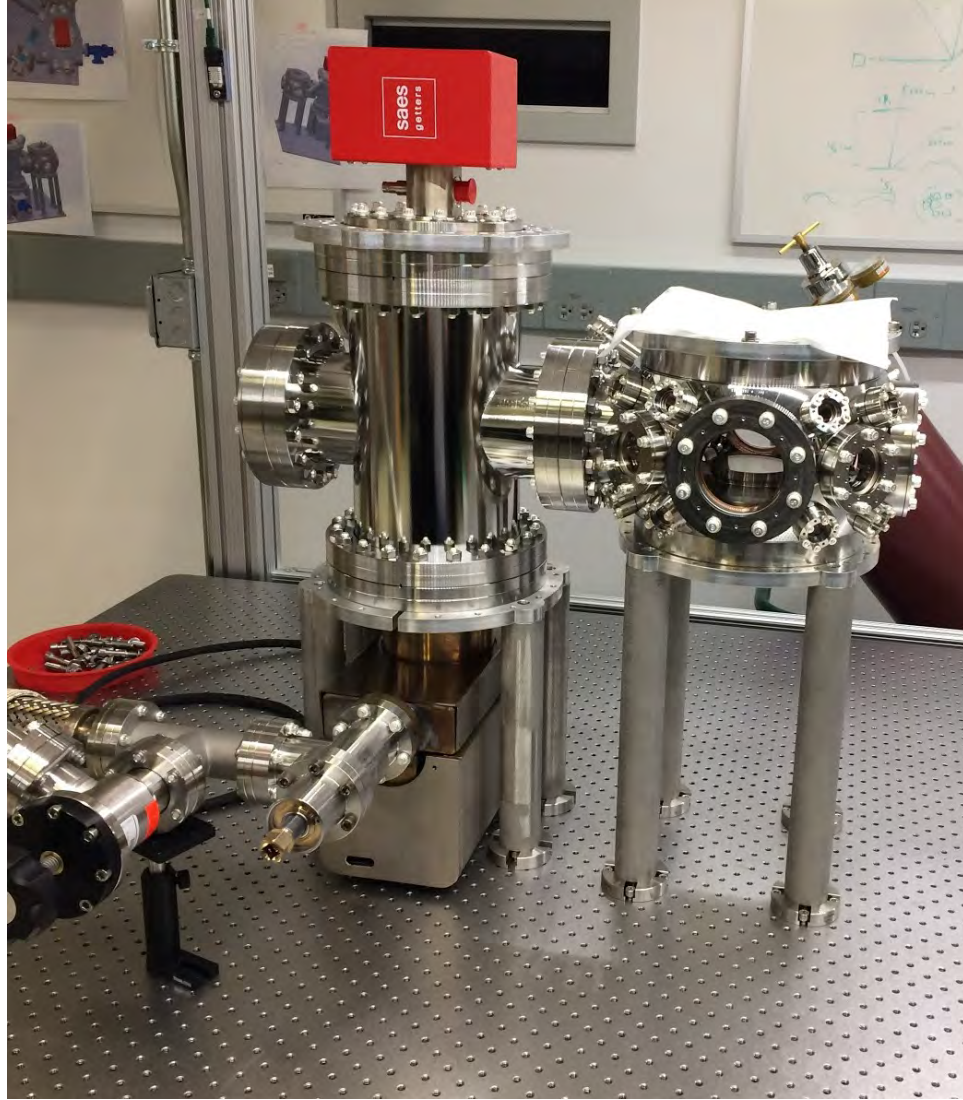


Figure 3.3: Photo of the chamber, before the source was attached or the breadboards were assembled.

The science chamber is based off a modified spherical square base from Kimball Physics. It has two recessed 2.69 in viewports on the top and bottom. These bucket windows hold our MOT bitter coils, and a high-resolution imaging system, and they

were designed in combination with the chamber so the mini viewports have line of sight to the atoms. There are three CF 4.5 in viewports, four CF 2.75 in viewports, and fifteen mini CF 1.33 in viewports. Two of the mini viewports opposite each other are specifically antireflection (AR) coated for the Rydberg laser wavelength (317 nm), this laser is mentioned in Sec. 3.3.6. Only one mini flange doesn't have a viewport and is instead used for the input of the source. The source is supported by 2 in diameter posts, and clamped in place with added rubber to avoid torque on the narrow connection to the main science chamber. There is a CF 6 in flange off of the science chamber going to a cross that holds a final CF 6 in viewport, a Nextorr-1000 non-evaporable getter (NEG) pump, and a 75 L/s ion pump with a port for a turbo pump. The turbo pump is only used for initially pumping down the chamber and is disconnected when we are operating the experiment. Our vacuum chamber has a pressure of $\sim 1 \times 10^{-11}$ Torr, which is low enough for any science direction we might take. Two custom breadboards are mounted around the chamber, one slightly below the center of the chamber so that standard Thorlabs optics are centered, and one breadboard above the chamber with cutouts for vertical optics.

3.1.4 Grating Chamber

The vacuum chamber for the grating MOT experiment was assembled rather quickly and had a few iterations before we arrived at our current configuration. The main part of the chamber is based on the design in Refs. [47, 57], and shown in Fig. 3.4. Most of the apparatus is constructed within a nitrogen-filled glove bag

to avoid oxidation of the strontium. The strontium source consists of a 3D printed titanium dispenser. The dispenser is a small cylinder, 5 mm in diameter, and 13 mm long, one end of the cylinder has a cap that can be opened to load it with a few mg of strontium, and closed tightly once fully assembled. The cylinder has a tab on both ends to connect to the current feedthroughs and a small slit to generate the atomic beam of strontium. This dispenser is generally operated at 12 A, and gets up to the operating temperature in a few minutes, but we have pushed this up to 14 A for short periods to increase the flux. Underneath the dispenser is a NEG capaciTorr Z100 pump, which is reactivated every time we open the chamber to air. Previous iterations of the apparatus had an atomic shutter, this was made from a 3D printed titanium piece and controlled through an actuator on a “micro” flange. The actuator turned an axle, which would rotate the titanium piece above the dispenser to block the atomic flux, or to the side to allow for flux. In practice, the shutter would consistently snag on pieces within the chamber, often getting stuck fully or halfway closed, and we would have to break vacuum to fix it. Additionally, the micro-flange is difficult to seal and the hardware often strips. We decided to scrap the shutter due to these issues, however, the flange remains as it is part of a custom piece, and thus we have to work around this flange when designing coils or magnet holders.

Above the dispenser is the differential pumping tube. The tube is 3 cm long and has a 3 mm diameter, which corresponds to an N_2 conductance of .11 L/s. This isolates the trapping region from the source. In addition, since the magnetic field extends into this region, the length of the tube acts as an effective Zeeman

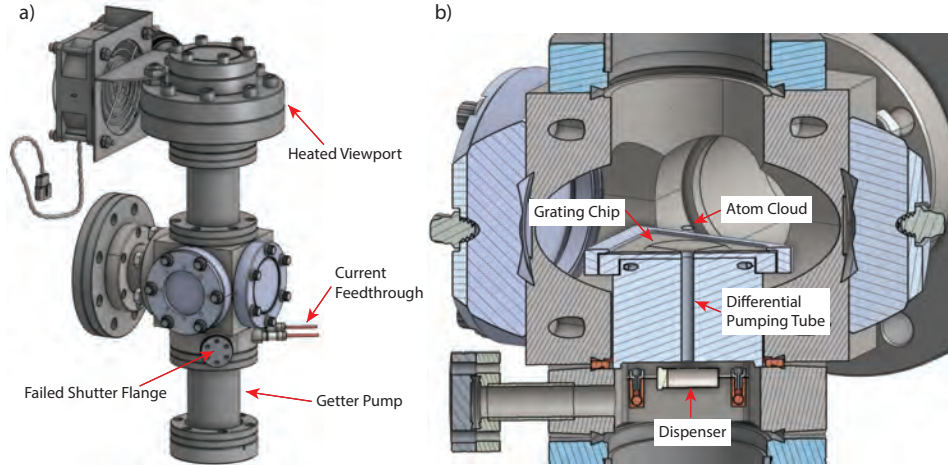


Figure 3.4: Renders of the final grating vacuum chamber. a) An external view of the chamber, the getter pump sits below the dispenser which is supported by the current feedthroughs. The vacuum tee, and connected ion and turbo pumps are not shown, nor are the coils. b) A cutaway to show the internals of the vacuum chamber. The dispenser sits below the differential pumping tube and grating mount. The atom cloud sits above the grating at the center of the trap.

slower [58?]. The machined part with the differential pumping tube also acts as a mount for the grating chip. This piece was machined out of aluminum in the original apparatus, however, we found that this created strong eddy currents, which made the field switching between the blue and red MOTs impossible. We had the piece re-machined out of titanium, which has a resistivity 15 times that of aluminum, and this solved the eddy current issue.

The grating is held in place with a clamping piece (also remade out of titanium). Details about the grating are described in detail in Ref. [58] and Sec. 3.4.3. Three 2.75 in viewports provide optical access to the grating chamber.

The top viewport was placed directly on top of the cube in the original apparatus, however, over time we noticed strontium build up on this viewport. We

replaced this viewport with a heated viewport (VHW-150-G), which is heated to ~ 300 °C. We have not seen strontium buildup on this new viewport, and it has not increased the pressure in our chamber significantly; however, we had to add a 3 in nipple to the top of the chamber to accommodate the larger radius.

One port of the grating chamber is connected to a vacuum tee which has a 75 L/s ion pump attached and a valve to connect a turbo pump. All together in the final iteration of the chamber, after multiple leaks and redesigns, the apparatus has a pressure of 1×10^{-10} Torr.

3.2 Magnetic Fields

Both the main apparatus and the grating MOT apparatus use approximately quadrupole fields to generate the magnetic field gradient for the trap, and shim coils to make fine adjustments to the field zero. We have used a variety of techniques to generate these magnetic fields in the laboratory, from simple coils, traditional water-cooled coils, bitter coils, and permanent magnets. This section will expand upon these different techniques.

3.2.1 Bitter Coils

One of the main changes to the new system is the introduction of bitter coils to replace the wound coils of the previous setup. Our bitter coils are based off of the design in Refs. [59, 60]. They consist of flat copper disks stacked on top of each other and are then water-cooled in parallel between the plates that make up the coil.

This allows for more efficient water cooling than running the water serially through the coil wire; thus bitter coils are often able to run at higher currents than wound coils. We want the ability to operate at high fields for possible future experimental directions, specifically implementing orbital Feshbach resonances, similar to Ref. [61].

Fig. 3.5 shows a diagram of our bitter coil design. Between each disk is an insulating polyester layer with channels cut into it to allow water to flow and a small copper section to allow current to flow between the disks. Each of our coils are formed by two concentric disk stacks, with a thicker rigid brass base to provide structure. Current flows from the bottom of the outer stack to the top, where a special copper layer connects the two stacks and allows the current to flow down the inner stack (depending on the polarity). On top of each coil is a stainless steel water manifold that distributes water through the various channels and defines the direction of flow.

The construction of the bitter coils revealed many issues with the design. Initially, the plan was to seal the coil closed using threaded rods, two of which would be brass and used as current inputs. In practice, the brass rods would easily break before the coil was sealed, so the final design includes copper bars as the current connection. Getting the coil to seal watertight also posed problems. The Teflon insulation alone was not enough to seal the coils, so we added rubber gaskets that outlined the channels in the Teflon. With the tight tolerances of our design, these gaskets could easily shift slightly out of place and spring a leak, making assembly quite difficult. In the end, we also insulated the entire coil with silicone rubber

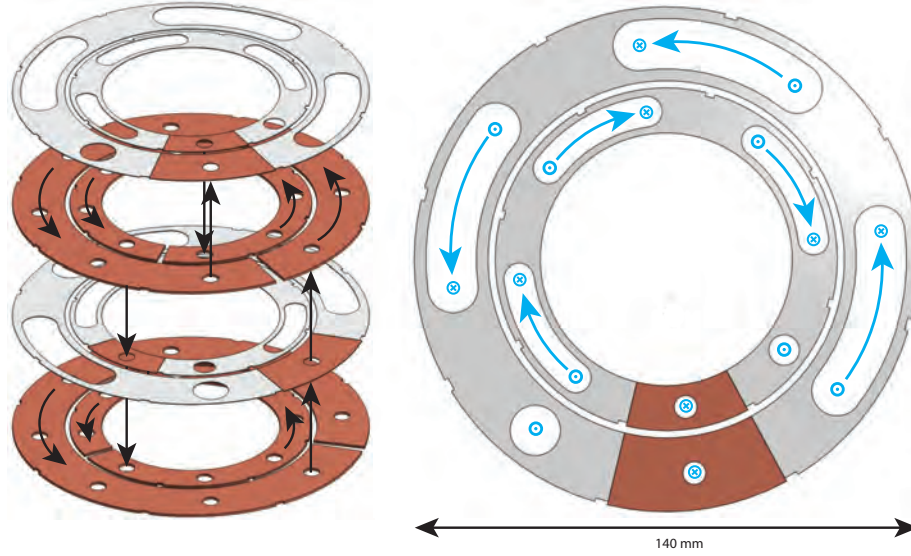


Figure 3.5: Exploded diagram of the bitter coil. The black lines show how current flows up and down the concentric coils. The blue lines show an example of how water flows through an insulating layer between the plates.

sealant to make sure it would not leak. Lastly, since the water distribution manifold is stainless steel there was a connection to ground through the optishield we put in the water. While we were able to increase the resistance to ground by using a non-conductive version of optishield for corrosion prevention, we also had to isolate the control electronics to prevent current from flowing through ground when the circuit was supposed to be closed. The coils have an inductance of $\sim 200 \mu\text{H}$ and give a field gradient per current of approximately 1 G/cmA axially when mounted on the chamber. Fig. 3.6 shows a photo of a completed bitter coil.

The control electronics for this system were designed by Ananya Sitaram and Alessandro Restelli. They use a metal–oxide–semiconductor field-effect transistor (MOSFET) bank with feedback from a current sensor with varistors for protection.



Figure 3.6: Photo of a completed bitter coil before we insulated it with silicone rubber sealant. Here you can see the total size of the coil and the water distribution manifold, the top flange is 8 in in diameter. The rubber gaskets provide insulation between the manifold and the sealing screws. The copper bars we use as current inputs, and the final silicon sealant are not shown.

The first iteration had an insulated-gate bipolar transistor (IGBT) to further decrease the switching time, the electronics for that section would break from time to time, and it made no difference in the MOT transfer efficiency, so we bypass it in the final design. Currently, the coils only operate in an approximately anti-Helmholtz configuration, but there is a break in the welding cables that supply the current for future integration of an h-bridge to operate in a approximately Helmholtz configuration.

3.2.2 Grating Coils and Magnets

We went through several configurations of coils and magnets for the grating MOT apparatus. The first iteration of the apparatus used purely permanent magnets and is described in Ref. [58]. These magnets were arranged somewhat asymmetrically to make room for the shutter flange but were still able to create a quadrupole field with an axial gradient of 62 G/cm. Unfortunately, this configuration is static and doesn't allow us to switch to the red MOT gradient.

The second configuration incorporated both magnets and coils. Half of the magnets were removed and 99 turn coils were wound above them. The coils supplied approximately 30 G/cm axially when fully on and were controlled with a bipolar MOSFET controller. This allowed us to operate the coils to add to the gradient of the permanent magnets and reach the blue MOT gradient, or switch the currents and operate them to cancel the permanent magnets and reach the low gradient of the red MOT. The bipolar coils with separate control also allowed us to use the same coils as the shim coils for the vertical direction. In practice, we ran into a few issues with these coils. We were running at the maximum voltages of our power supplies and were seemingly still not able to fully cancel the gradient of the magnets. However, this was most likely due to the eddy current in the differential pumping tube rather than an issue with the coils themselves. Furthermore, these coils were not water-cooled, so they would overheat on the order of 30 s. While this could be mediated by simply adding 1-2 s to the duty cycle of the experimental shots, it made continuous alignment and optimization of the blue MOT difficult. We decided

to scrap these coils in favor of more traditional water-cooled coils, but I think this configuration would work in the future given that we have solved the issue with the eddy current.

The final configuration has two sets of coils: one set of large water-cooled coils to supply the blue MOT gradient and one small set of coils to supply the red MOT gradient. The large coils are very similar to the coils from the apparatus in the old laboratory and supply approximately 0.55 G/cmA axially. We operate these coils at 100 A for the blue MOT, to get an axial gradient of 55 G/cm. Unfortunately, our bipolar controllers cannot operate at this current, so we use an IGBT to completely and quickly turn these coils off and use a separate, more carefully controlled, set of coils for the red MOT. This smaller set of coils is not water-cooled and gives us 0.2 G/cmA. We operate these coils with the bipolar controllers allowing us to use them as the vertical shim as well. Near the max allowable current of the controllers, they give us the 4 G/cm axial gradient we use for the red MOT. We would like these coils to be a bit stronger, but cannot wind any more turns due to the current geometric constraints of the chamber, a limited redesign would probably be able to improve on this.

3.2.3 Shim Coils

Both the main and grating apparatus use shim coils to move the zero of the MOT field gradient or zero the field in a dipole trap. There are shims along all three axes on the main chamber, six coils in total. The vertical coils have 78 turns each,

and the horizontal coils have 60 each, except for the coil on the side of the chamber that connects to the pump cross. Since this coil is wrapped around a larger flange than its partner, it has 31 turns to compensate for the larger size and maintain a flat magnetic field profile in the center of the chamber. All three sets of coils are run with computer-controlled power supplies.

The grating MOT chamber only has shim coils in the transverse directions; the axial direction is shimmed with the MOT coils. Three of these coils are wound around 3D printed cage mounts and have 30 turns each. We have computer control only in the y direction (along the imaging axis), while the x shim is manually tuned. In both systems the shims are not powerful enough to have much effect on the magnetic field gradient of the blue MOT but can move the zero of the red MOT field gradient significantly.

3.3 Laser Systems

Working with strontium requires many different laser systems. In the laboratory, we have one optical table for the blue lasers and dipole lasers, and another for all the red lasers, including the repump lasers. The last optical table is reserved for the vacuum systems and a few miscellaneous lasers that will only be discussed in brief. This section will be a relatively brief overview of these laser systems, as many of the lasers were transferred without alteration from the old lab and are described and diagrammed in Refs. [5, 55, 56].

3.3.1 Blue 461 nm System

The first cooling stage of strontium, the blue MOT, uses light at 461 nm. Our 461 nm system consists of two lasers. The “spectroscopy” laser is a Toptica DL pro, which we lock to a hollow cathode lamp with polarization rotation spectroscopy [62] on the $^1S_0 \rightarrow ^1P_1$ transition in ^{88}Sr . This lock is sensitive to power variations and magnetic field, so we stabilize the power of the spectroscopy beams, and cover the hollow cathode lamp with mu-metal to provide magnetic shielding. The old setup had an optical isolator (from a different laser system) nearby the cell, the field from that isolator was written onto the mu-metal so when it was removed this field made the lock unusable, and the mu-metal had to be replaced. The power lock is even more important, as small shifts in power will significantly move the frequency, and it only works well when the maximum power is significantly higher than the lock point. Even with the magnetic shielding and power lock, the frequency of the lock can shift by a few MHz over an hour or two. The laser is locked 100 MHz below resonance using an acousto-optic modulator (AOM) that shifts the spectroscopy laser beam up in frequency.

The main cooling laser is a new TA-SHG pro, which can output ~ 900 mW of light. Currently, this light is split into four main paths with polarizing beamsplitter (PBS) cubes, light for the 2D MOTs, the Zeeman slower, the main MOT beams, and the imaging probe beam with approximately 20 mW, 25 mW, 9 mW per beam, and 300 μW out of the respective fibers. These powers can easily be adjusted with $\lambda/2$ waveplates before each cube. The light for the 3D MOT is coupled into a 1 to

3 fiber splitter, and power locked using the attenuation of that path's AOM, and a photodiode on the experiment table. Fig. 3.7 is a diagram of this system.

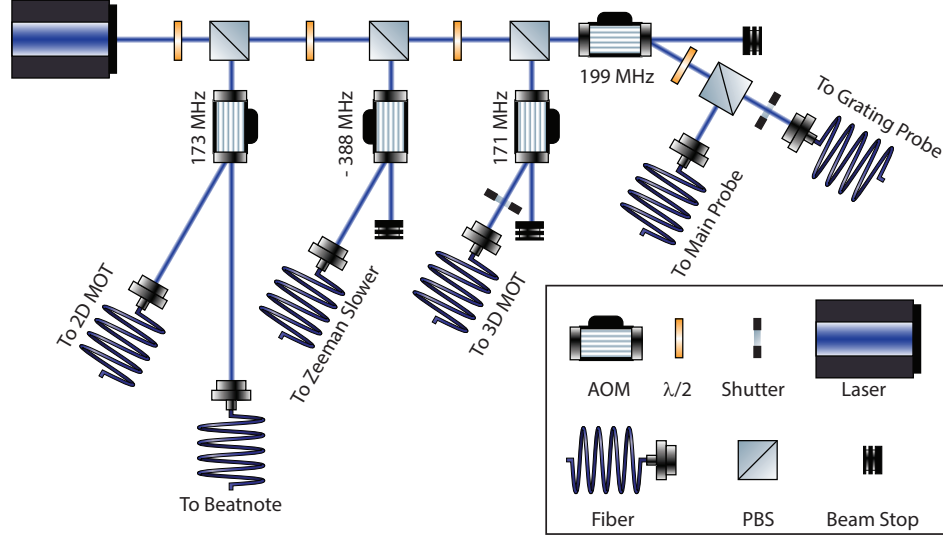


Figure 3.7: Diagram of the blue laser system. The mirrors are omitted for simplicity. Each arm of the blue laser system has its own AOM, and waveplate to control the power. Optics components from Ref. [63].

The unshifted light from the zeroth order of the 2D MOT AOM is used to lock the cooling laser below the spectroscopy laser with a beatnote lock. This beatnote lock is somewhat noisy, and while it doesn't seem to significantly effect the loading of the blue MOT it can make fluorescence imaging less consistent. The beatnote is ~ 100 MHz for ^{88}Sr and increases according to the isotope shifts for the other isotopes, listed in Tab. 1.3, so that the unshifted laser is locked ~ 200 MHz below resonance. An individual AOM shifts each arm of the system to its respective frequency relative to the lasing frequency: +160 MHz for the 2D MOT beam, -388 MHz for the Zeeman slower beam, +155 MHz for the main MOT beams, and

+200 MHz for the probe beam. This setup is changed and simplified from the old apparatus, in particular, the frequencies have been shifted to remove a double pass AOM, and decrease the amount of wasted power.

For the final iteration of the grating MOT system we use this laser for its high power. When running the grating MOT the fiber splitter is replaced with a single fiber going to the grating MOT beam input, and almost all the power is put into this arm giving ~ 200 mW after the fiber. Some power is still reserved for the beatnote lock, and probe beam. We can switch between the main apparatus probe, and the grating MOT probe using a waveplate and cube.

3.3.2 Repump Lasers

Our apparatus uses two different repump frequencies as mentioned in Sec. 2.1.2. We have two homemade, extended cavity diode lasers (ECDLs) in the Littman-Metcalf configuration to provide light at these frequencies. Both these lasers are locked through slow feedback using the reading from the wavemeter (High Finesse, WS7). We have a Python program that reads the wavemeter and provides this feedback. Yanda Geng remade this program based on the old Labview code to operate more quickly and fully separate our lab from Labview. The approximate locking frequencies are shown in Tab. 3.1. The wavemeter has a precision down to 1 MHz, but can drift by a few MHz over a day.

The 679 nm laser is sent through an AOM to provide more precise control of timing and future imaging out of the $^3\text{P}_0$ state. The atom number is more sensitive

Repump Frequencies (THz)	679 nm	707 nm
Bosons	441.3326	423.9136
Fermion	441.3301	423.9130

Table 3.1: The wavemeter values of the repump laser frequencies for the bosons and the fermion. Note that the wavemeter drifts day to day, and the 679 nm laser is shifted by 207 MHz above the wavemeter measurement by an AOM.

to the 707 nm repump which has to be broadened to cover the hyperfine transitions of ^{87}Sr , as mentioned in Sec. 2.1.2. For both of these reasons, we need a larger amount of power in this repumper. To achieve this we injection lock another diode laser seeded with the homemade ECDL to provide sufficient power. We add sidebands to the output of this laser with a fiber electro-optic modulator (EOM), which is driven at many different frequencies, as discussed in Sec. 4.2, when we are trapping ^{87}Sr .

3.3.3 Red 689 nm System

Operating the red MOT requires a laser at 689 nm with a narrow linewidth to address the narrow transition. This narrow linewidth is achieved by locking a laser to an Ultra Low Expansion (ULE) glass cavity from Advanced Thin Films (ATF-6010-4) with a cavity linewidth of ~ 6 kHz, I will refer to this laser as the “cavity” laser. The cavity laser is sent through a double pass ~ 1 GHz AOM. The double passed diffracted order is sent through a fiber to a free-space EOM to lock the laser to the cavity with a Pound-Drever-Hall (PDH) lock. This cavity and lock are described in detail in Ref. [55]. The free spectral range (FSR) of the cavity is approximately 1.5 GHz, so together with the double passed AOM, this allows us to have full flexibility in the locking frequency of the cavity laser. The zeroth order

of the AOM is sent to a beatnote lock with an ECDL laser. We call this ECDL laser the “long Steck” since it is based on the design by Dan Steck [64] (and is a bit longer than a similar laser in the laboratory we call the “short Steck”). In the previous system, this beatnote was measured with a large network of fiber splitters that picked off light for the wavemeter and measured beatnotes for a variety of lasers that are no longer in use. We replaced this setup with a single fiber splitter to conserve power and simplify the system. Additionally, we added computer control to the beatnote frequency, a new feature of the current apparatus.

The long Steck is used as the main cooling laser for the bosonic isotopes. The beam path to the experiment has an AOM and shutter to shift, shutter, broaden, and control the power of the beam. Then the light is sent through a 1 to 4 fiber splitter to provide the light for the MOT beams. Typically we have $\sim 600 \mu\text{W}$ per beam for the red MOT.

As always, for the fermion, there are additional complications. In this case, the light from the long Steck is used as the stirring laser, mentioned in Sec. 2.1.3, as opposed to the cooling laser. Some of the light from the cavity laser is sent to inject another diode laser which we then use for the cooling and can provide more power, $\sim 2 \text{ mW}$ per beam.

For the grating MOT, we replace the fiber splitter, which splits the main red MOT beams, with a fiber that seeds a Toptica BoostA tapered amplifier (TA). The fiber-coupled output of this TA supplies $\sim 90 \text{ mW}$ of power when operating at its maximum current, and follows the broadening of the red laser well. While searching for the fermionic red MOT in the grating system we have used a few

different configurations informed by our simulations. We tried seeding the TA with the injection laser, and using it as the cooling laser, but with opposite circular polarization. We also tried seeding the TA with the long Steck as the cooling laser and operating the injection on the $F = 9/2 \rightarrow 7/2$ transition for stirring with the same circular polarization for both. Lastly, we are trying the same configuration but seeding the TA with both the injection, and long Steck combined onto a fiber splitter.

3.3.4 Clock Laser System

We need another narrow laser, operating at 698 nm to address the clock transition. This laser is another long Steck style laser, which has one beam path with an AOM going to the experiment and another which is sent through a fiber EOM for a PDH lock to another ULE cavity. This second cavity is more carefully engineered and sits on a vibration isolation stage and inside a sound-insulated box. A very detailed description of the cavity is presented in Ref. [5], although a new box had to be made for the new laboratory. The fiber EOM is driven with two different frequencies one at 18 MHz for the PDH lock, the other at ~ 800 MHz to add sidebands to the laser. We lock one of these ~ 800 MHz sidebands to the cavity, rather than the carrier, This gives us the flexibility to lock the frequency of the laser away from a cavity resonance. Previous iterations of the clock laser system had a pre-stabilization cavity, to initially narrow the laser before the ULE cavity, but this was scrapped as it didn't improve the lock. The beatnote between the red cavity

laser and the clock laser is ~ 200 Hz wide, putting an upper limit on the linewidth of the clock laser.

3.3.5 Dipole Trap Laser Systems

We have two dipole trap lasers. The first is a 1064 nm, 30 W laser that was transferred directly from the old laboratory. This laser system had four beam paths for different dipole traps, but we have only used two of them in the new laboratory. Each path has an AOM and mechanical shutter, for quick and complete extinction respectively. We send ~ 2 W to the science chamber for each dipole trap, but this power can be lowered once the traps have been aligned.

The second dipole trap laser is an M squared SolsTiS ti-sapphire laser. This laser has a wide frequency range and can provide ~ 4 W at 813 nm, the magic wavelength for the clock transition. We had a lot of issues with this laser in the past. Previously we pumped the laser with a 532 nm IPG laser, but multiple of these lasers failed, and eventually, the ti-sapphire laser failed as well. We now have a new system with a pump from M squared which we hope will be more stable and are working on getting the magic wavelength dipole traps working.

3.3.6 Miscellaneous Lasers

The laboratory has a few other lasers that I would like to mention here briefly. First is the old blue laser system from the old apparatus. It is another TA-SHG Pro from Toptica, similar to our current system, but is nearing the end of its life,

and can only put out $\sim 300 - 400$ mW. We previously used this system to run the grating MOT, but after finding that we were power limited we switched to the new laser system instead. The current plan is to use this system for blue detuned dipole traps, specifically to make the walls of a box trap, discussed in Sec. 6.1.

The second laser is a Toptica TA-FHG Pro, this is a fourth harmonic generation laser at 317 nm laser to excite Rydberg levels off of the clock state, as discussed in Sec. 6.2. This laser sits on the table with the vacuum chamber due to the difficulty of efficiently fiber coupling ultra-violet light. Currently, this experiment has been put on hold, so the laser isn't currently in use.

The final laser I would like to discuss is the frequency comb in the lab. The comb is from Menlo Systems (FC1500-250-WG) and is a resource for the JQI as a whole. The comb acts as an absolute frequency reference by translating optical frequencies to a combination of three radio frequencies: the repetition rate (~ 250 MHz), the offset frequency (~ 35 MHz), and the beatnote between the closest mode of the comb and the laser being measured. The comb was a particularly valuable resource for the work in Ref. [5], but is also useful as a way to get lasers close to the narrow transitions in strontium. The comb has amplifiers to provide power centered at 688 nm, 698 nm, 780 nm, 1112 nm, and 1166 nm, although the 688 nm and 698 nm amplifiers can't be operated simultaneously. In addition, there is a photonic crystal fiber to broaden the laser and provide power over a wide range of wavelengths from 640-900 nm.

3.4 Experiment Side Optics

3.4.1 Main Experiment Optics

So far the main experiment requires optical access for blue and red MOT beams, two dipole trap beams, and two imaging paths, with the future addition of a blue detuned dipole trap. Figs. 3.8 and 3.9 show the optics on the vacuum chamber that are used to shape and combine these beams.

The blue MOT beams are collimated and their polarization is set and cleaned up with waveplates and PBS cubes. One of the blue MOT beams has a pickoff for power stabilization. The red MOT beams have a similar setup and are combined with the blue beams on a dichroic mirror. The horizontal blue and red MOT beams enter the chamber at 45° relative to the imaging axis and are rotated to circular polarization with dual-color $\lambda/4$ waveplates. These waveplates are actually $\lambda/4$ for 461 nm and $3\lambda/4$ for 689 nm, so the red beams have an extra $\lambda/2$ waveplate to correctly account for this. There are mirrors to steer the blue and red MOT beams separately, as well as to steer them together. The blue and red beams have $1/e^2$ power distribution radii of ~ 6 mm and ~ 3 mm respectively giving max I/I_{sat} of $\sim .4$ for the blue beams. We previously placed a telescope in these beams to further expand them, but we found that we did better by taking it out, however, it's likely the atom numbers could be improved by further expanding these beams. Simple retro-reflecting mirrors create the opposing beams, with waveplates to maintain the correct circular polarization. The repump fibers are mounted to the chamber with

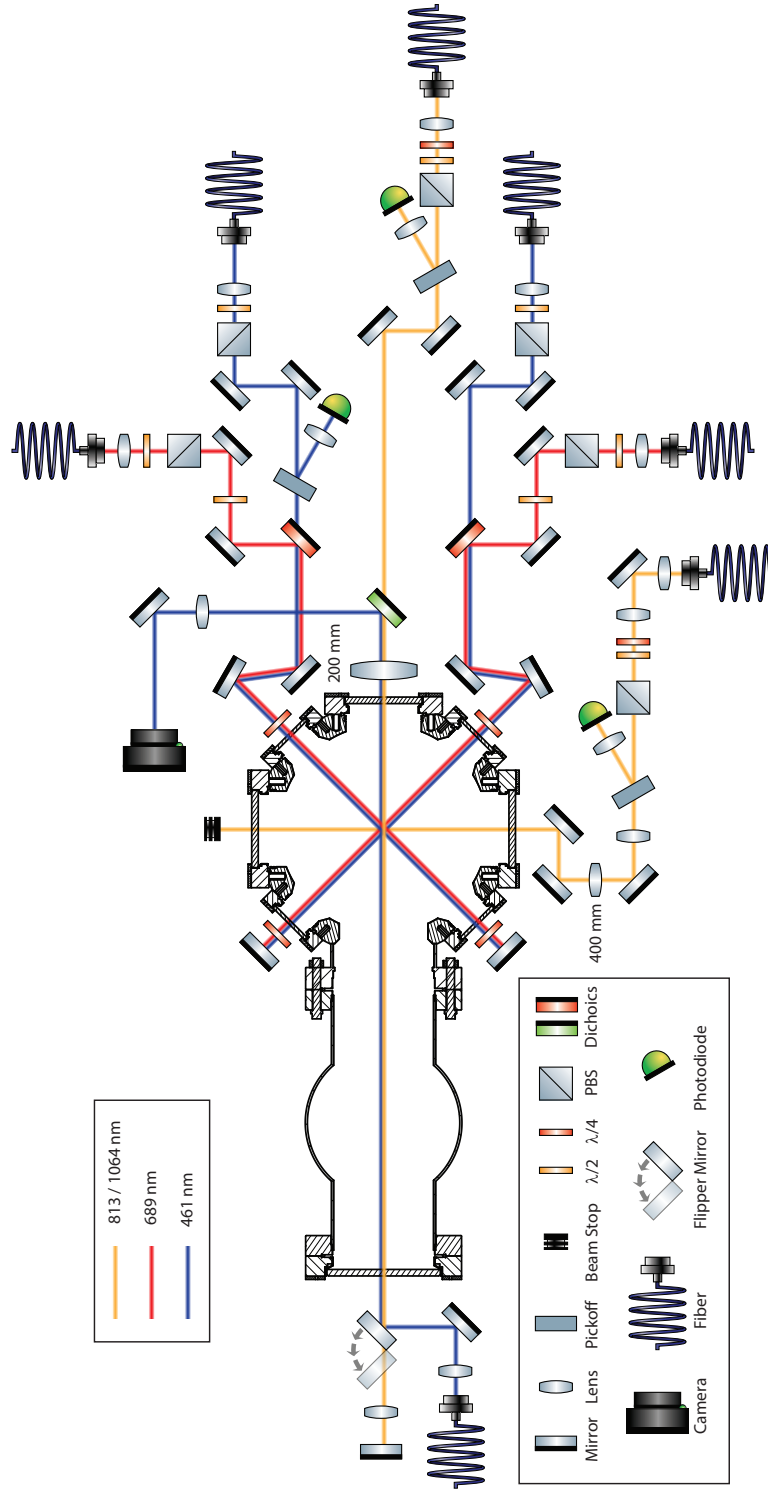


Figure 3.8: The horizontal optics for the main experiment. This diagram is not to scale. Optics components from Ref. [63].

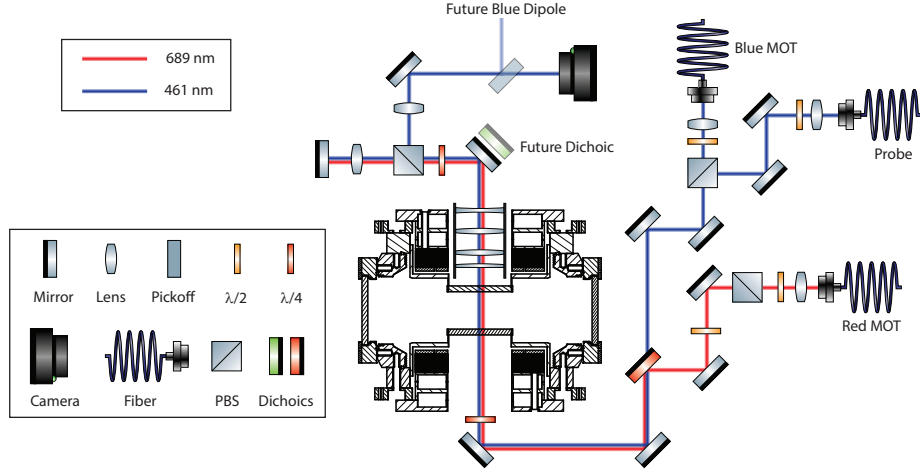


Figure 3.9: The vertical optics for the main experiment. The vertical beam paths have the high resolution imaging stack. The 50% opacity dichroic and pickoff are for the future implementation of vertical red and blue detuned dipole traps respectively. This diagram is not to scale. Optics components from Ref. [63].

Thorlabs cage mounts, and viewport washers and adapters machined at the JQI, and enter on two of the mini viewports.

The x -direction dipole trap simply enters the chamber perpendicular to the horizontal imaging, this beam is focused onto the atoms with a 400 mm focal length lens, and the beam is shaped to provide a $\sim 30 \mu\text{m}$ waist. The cross dipole trap counter-propagates with the horizontal imaging probe beam. It is combined with this beam on a dichroic mirror and is focused onto the atoms to create a $\sim 30 \mu\text{m}$ trap with a 200 mm focal length lens, which is also used as the first lens of the imaging system. On the other side of the chamber, there is a flipper mirror to switch between imaging and retro-reflection of the dipole beam to create a lattice. Both dipole traps have pickoffs to photodiodes for power stabilization. The horizontal

imaging system is a simple two-lens system with a magnification of 1.4.

When the data in Ch. 4 was taken the vertical beam paths looked much like the horizontal MOT paths. Since then we have added optics to allow for high-resolution imaging and the box dipole traps. Parts of this system are still under construction. The MOT beams are launched from below the chamber, along with an added probe beam. On top of the chamber a high-resolution imaging stack, described in Sec. 3.4.2, sits inside the bitter coil. We plan to integrate a dichroic mirror for the implementation of future 813 nm vertical dipole traps, however, we have had some issues with fringing from this dichroic, so for now it is replaced with a simple mirror. After a dual-color $\lambda/4$ waveplate a PBS beam splitter separates the MOT beams from the probe beam. The MOT beams are collimated and retro-reflected, while the probe beam is sent to the “PIXIS” (PIXIS Excelon 1024B) camera from Princeton Instruments. Future blue detuned dipole beams will be combined with the probe beam with a simple beam splitter.

The AOSense source has beam shaping optics for the 2D MOT and Zeeman slower laser beams that hang underneath the top breadboard. The 2D MOT beam is shaped into an elliptical beam with cylindrical lenses, before being split between the two 2D MOT inputs. The Zeeman slower beam is circularly polarized and sent through the heated viewport. Fig. 3.10 shows these optics.

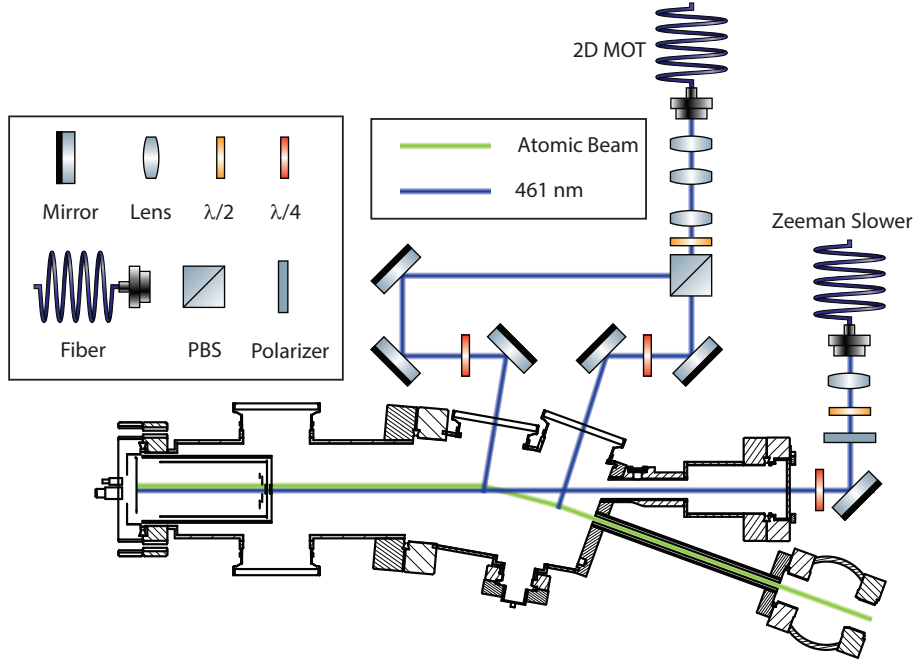


Figure 3.10: The optics for the AOSense source. The atomic beam is also shown including the progressive angles of the 2D MOTs. The internal optics are not shown as they are proprietary to AOSense. This diagram is not to scale. Optics components from Ref. [63].

3.4.2 Imaging Stack

For the new strontium apparatus, we designed a high-resolution imaging lens stack. The design for the stack was inspired by Ref. [65]. This design aims to create a four lens stack out of stock Thorlabs lenses. We start by putting four lenses and the viewport of the vacuum chamber into Zeemax and optimizing the imaging resolution for 461 nm using the surfaces of the lenses. Then we replace one of the lenses with the closest Thorlabs part and redo the optimization. This process is iterated until all the lenses and the spacings are fixed and optimal. Fig. 3.11 shows

the lens stack and spacings.

The final lens stack is assembled with aluminum spacers that were precisely machined to match the spacings. Since this lens stack is fitting inside one of the bitter coils we cut slits into all the spacers, the lens tube, and even the retaining rings, and filled the slits with Teflon to avoid eddy currents in these parts. The stack has an NA of 0.29, and a resolution of $\sim 1 \mu\text{m}$ at 461 nm measured on a test setup.

3.4.3 Grating Optics

The optics on the grating MOT apparatus are less than ideal; space is limited and they are supported by long posts rather than a proper breadboard. Fig. 3.12 shows a diagram of this setup. The blue MOT beam polarization is cleaned up with a cube, this cube is useful for alignment as the back-reflected zeroth order of the grating will be rejected by the cube. The beams are reflected by a dichroic mirror before the circular polarization is set with a dual-color waveplate, so that the dichroic doesn't scramble the polarization. A telescope then expands the beam before it enters the chamber and is diffracted by the grating. The red MOT beam has a similar setup, with extra waveplates to provide separate control of the red and blue polarization. The blue MOT beam has a ~ 9 mm radius and the red MOT beam has a ~ 5 mm radius. In the future, it could be possible to use planar optics to simplify and shrink the MOT beam shaping system[66].

The grating chip, shown in Fig. 2.2, has three grating sections arranged in

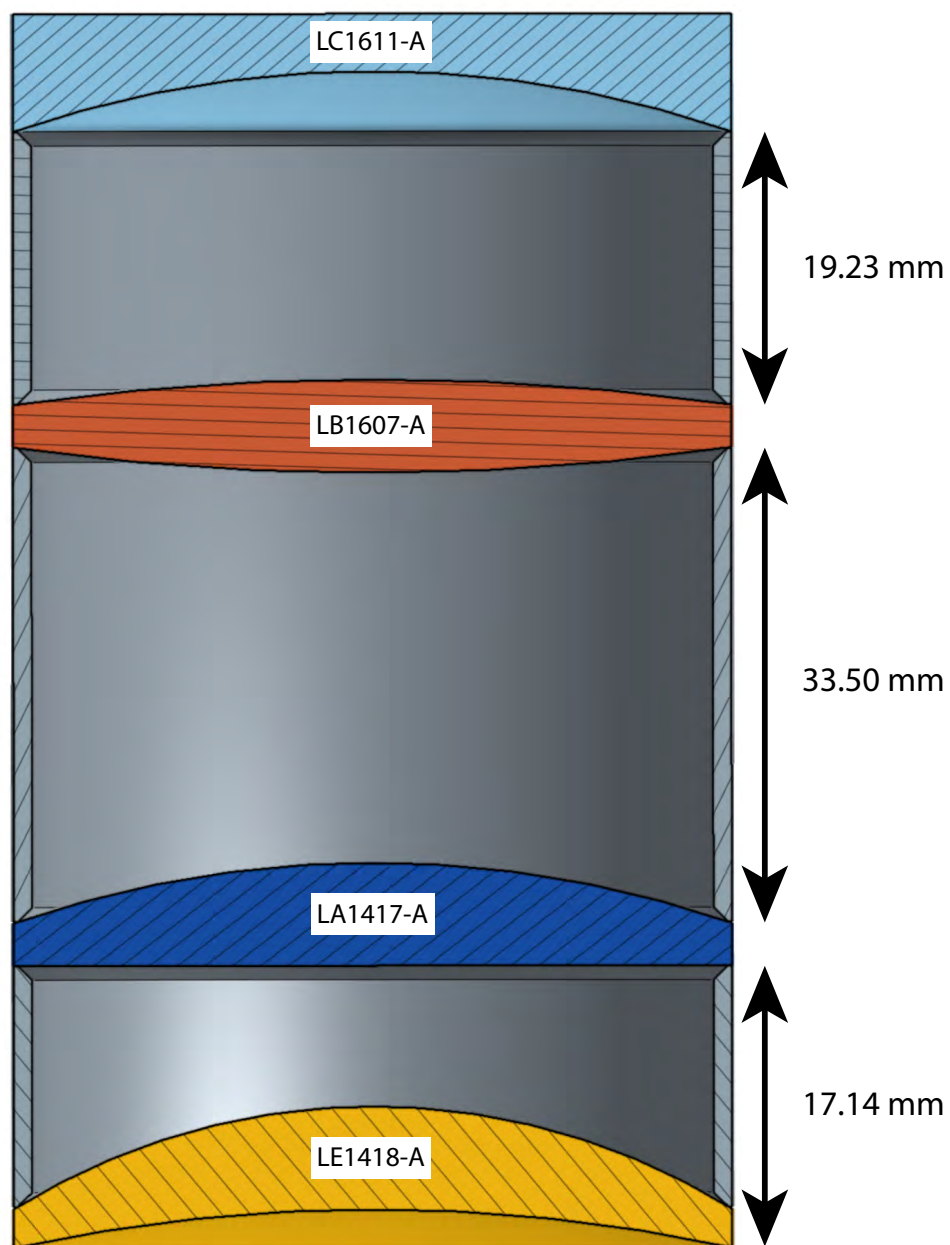


Figure 3.11: The imaging stack we designed for the new strontium apparatus. The Thorlabs part numbers for the lenses, and the spacings between them are shown.

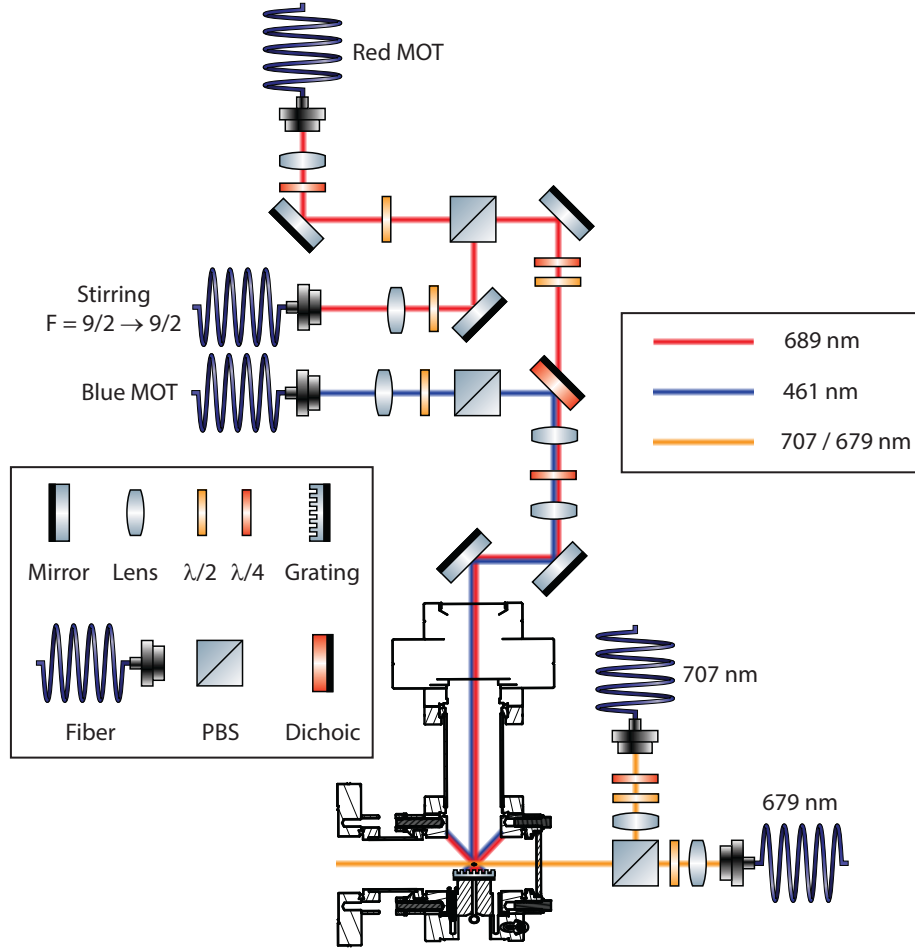


Figure 3.12: The optics on the grating MOT. The stirring laser for operating on the $F = 9/2 \rightarrow 9/2$ transition has its own input. The repumps come in on the viewport opposite the pump tee. The probe and imaging system are not shown and have a simple linear setup. This diagram is not to scale. Optics components from Ref. [63].

a circle with an inch diameter. The center of the chip has a triangular hole that sits over the differential pumping tube. Each grating section diffracts 32% of the blue light into each of the ± 1 orders, at an angle of 27° relative to the normal, and $\sim 37\%$ of the red light at an angle of 43° . With these angles, there is enough overlap (43%) between the blue and red capture regions for the transfer between the MOTs to work.

The repump beams are combined on a PBS and sent through the viewport opposite the pump tee. The probe beam enters the chamber perpendicular to the input MOT beams and the repumps, and is imaged with a simple two-lens imaging system with a magnification of 0.25.

3.5 Computer Control

The old laboratory used SetList, a program created with labVIEW, to provide computer control for the apparatus. Unfortunately, the user base of SetList is becoming less active, and the user experience was a bit clunky in the old laboratory. In addition, we wanted to move away from labVIEW in general, so we decided to change our computer control scheme for the new laboratory. We transitioned to the labscrip suite of software [67], which had already been in use by Ian Spielman’s groups at the JQI. This software has a more active user base than setlist and is python-based allowing for easy extensibility. The software suite consists of four main programs: *Runmanager* defines the variables for experimental sequences, *Blacs* communicates with the hardware and runs each shot, and also allows you to

manually control all the hardware when a shot isn't running, *Lyse* allows you to write custom scripts to analyze data from each experimental run as they come in, and using *Runviewer* you can view the output of each experimental sequence. The map of connections between all the devices in the lab is encoded in a python file called the "connection table", and the logic of the experiment itself is written in another python file using the *labscript* library. All the data for each experimental shot is stored in an hdf5 file. This hdf5 file includes all the parameters of the shot, the experimental logic, any results (including images), as well as the results of your analysis scripts. The file is created with *Runmanager*, sent to *Blacs* where the shot is run on all the devices, and the results are written to it, and finally, it is sent *Lyse* for analysis. The organizational scheme of this new software keeps everything organized and in one place, while simultaneously allowing us to split experimental control over multiple computers. In addition, there are tools to automatically run optimization algorithms using the M-Loop python package. While these algorithms rely on a stable system, they have been useful in searching the large parameter space of the red MOT.

Labscript is built off of the use of programmable "pseudoclocks", with one "master pseudoclock" to provide the timing for each experimental shot. previously we used a pulseblaster for our master pseudoclock, but Ananya Sitaram and Alessandro Restelli designed an alternative, called Jane with more channels, and built-in digital buffers [68]. This master pseudoclock is used directly for TTL signals to a variety of devices, in addition to clocking "intermediate devices" such as National Instruments (now NI) cards, and Arduinos. The laboratory uses three NI cards (NI

PCI-6733, NI PCIe-6361) that have both digital and analog channels.

The AOMs are controlled in a variety of ways, their drivers, designed and built at the JQI, have: a digital input that turns the radio frequency (RF) power on and off, an analog input that controls the attenuation of the RF power, an analog input to control the frequency of the voltage controlled oscillator (VCO) inside the driver, and a TTL to control whether the frequency is determined by the VCO or an external RF signal. Some of these inputs are unused by some AOM drivers, and unfortunately, depending on the version of the driver, the RF TTL or attenuation might have the opposite sign.

The laboratory also has a variety of direct digital synthesizers (DDSs) and phase locked loops (PLLs) that are controlled using Arduinos. These are used for beatnote locks, and to control the broadening of the 707 nm repump laser. Each Arduino takes a clock TTL and a serial connection that sends it commands. The code on these Arduinos was originally written to work with SetList, but we were able to get them to work nicely with labscrip with a bit of modification and a new labscrip device script. They are capable of controlling multiple DDSs by setting the frequency and doing frequency ramps that are clocked internally by the DDS chip to avoid any lag from the Arduino. Often when controlling beatnote locks the frequency that the DDS is set to is multiplied by a PPL before the beatnote comparison happens, so this multiplication has to be factored in when programming an experimental shot. There are still a few DDSs which we have not yet implemented into labscrip, such as the DDSs that control the clock fiber EOM sidebands, or the ~ 1 GHz AOM for the red cavity laser, but future implementation of computer

control for these DDSs could be useful.

There are three cameras we use in the lab that are controlled by labscript. The “Flea” (FL3-U3-13Y3M) currently images the grating MOT, and the “Grasshopper” (GS3-U3-28S4M) currently provides the horizontal imaging, both of these cameras are from Point Grey, and work with NI drivers. The last camera is the PIXIS, which we are implementing into the vertical imaging system, and has its own software and driver. Labscript has a very nice camera class for cameras that integrate with NI software. The class has a widget in *blacs* that allows for manual continuous imaging, which has been a valuable resource for our optimizations of the blue MOT. The connection table defines separate attributes for the cameras in manual and buffered (during a shot) operation. Unfortunately, we had to write a new labscript class for the PIXIS, and so far have not been able to get continuous imaging to work.

To broaden the red cooling and stirring lasers, as well as perform SWAP, we use a Stanford Research Systems (SRS) function generator (DS345). The function generator has a TTL to turn the output on and off, and an analog input to control the amplitude of the oscillation. We switch between sawtooth and triangle waves manually. This signal is added to analog signals for the VCOs of the cooling and stirring laser’s AOM drivers in order to broaden the lasers.

The remaining digital channels control: laser shutters, camera triggers, the IGBT in the grating MOT coil circuit, turning on and off the grating red MOT coils current lock, and turning on and off the bitter coil current lock. The remaining analog channels control: shim coil currents, the grating blue MOT coil current, the set point for the top and bottom current of the grating red MOT coils, the set point

for the bitter coil current, and the set point for the blue laser and dipole beam power locks.

In conclusion, we have built a new and improved apparatus for ultracold strontium experiments. Our new apparatus has a high flux source, good optical access, and coils capable of providing high magnetic fields. We have eliminated the issue of strontium coating viewport windows from the old apparatus, and significantly decreased the size. Our old and new laser systems are fully integrated into the experiment and everything is managed through a more streamlined, and extendable computer control system.

Chapter 4: Validations

This chapter walks through the process of trapping strontium atoms in the various traps, the techniques we used to initially find appropriate trap parameters, and the result of each cooling and trapping stage. These traps include the blue and red MOTs of both the main apparatus and the grating system, as well as the dipole traps of the main apparatus. In addition I will discuss the imaging techniques we use to measure the trap properties. The experimental parameters and results presented here are typical values found empirically, but they can change significantly day to day, or depending on what we are specifically doing in the laboratory. I will use “field” in this chapter to refer to the magnetic field unless otherwise stated.

4.1 Imaging

We extract all the information we need from our trapped atoms through imaging. From measuring atom number, temperature, and peak optical depth for optimization, to techniques for finding initial parameters that will trap any atoms at all, imaging is fundamental. In our laboratory we use two simple forms of imaging: fluorescence imaging, and absorption imaging.

Fluorescence imaging is the simpler of the two; a probe beam of resonant light

interacts with the atoms, and the resulting fluorescence is captured, off the probe beam axis, by an imaging system, and focused onto a camera. In practice it is difficult to extract accurate atom numbers from fluorescence imaging, as the light collected depends on the solid angle of the optical system, and how that solid angle overlaps with the dipole radiation pattern. We mostly use fluorescence imaging for initial alignment and optimization of our blue MOTs, as fluorescence images can be taken continuously while the MOT is active. However, in our system, the blue MOT beams are acting as our probe for fluorescence imaging, as we do not have a dedicated fluorescence probe beam. This is not ideal as the MOT beams are detuned from the resonance. Thus, optimization of the MOT detuning using fluorescence imaging will not be accurate, as detuning closer to resonance will increase the fluorescence even with the same number of atoms. Despite this, in the grating MOT system fluorescence imaging has been a useful tool, as it has a better signal to noise than the absorption imaging in our apparatus. More details on how we image our clouds in the grating MOT system are described in Sec. [4.5](#) and [4.6](#).

Absorption imaging is the more accurate and versatile tool. To take an absorption image a low power resonant probe beam is sent along the axis of the imaging system where it is absorbed by the atoms. The light that is not absorbed is imaged onto a camera. This first image (I_{atoms}) is proportional to the transmitted intensity. Then a second image (I_{probe}) of the probe beam is taken once the atoms have fully left the interaction region. This allows you to account for the probe beam power, and spatial variations in the beam profile. We also take a background image (I_{bg}) with the probe beam off. Consider the amount of light dI that will be absorbed by

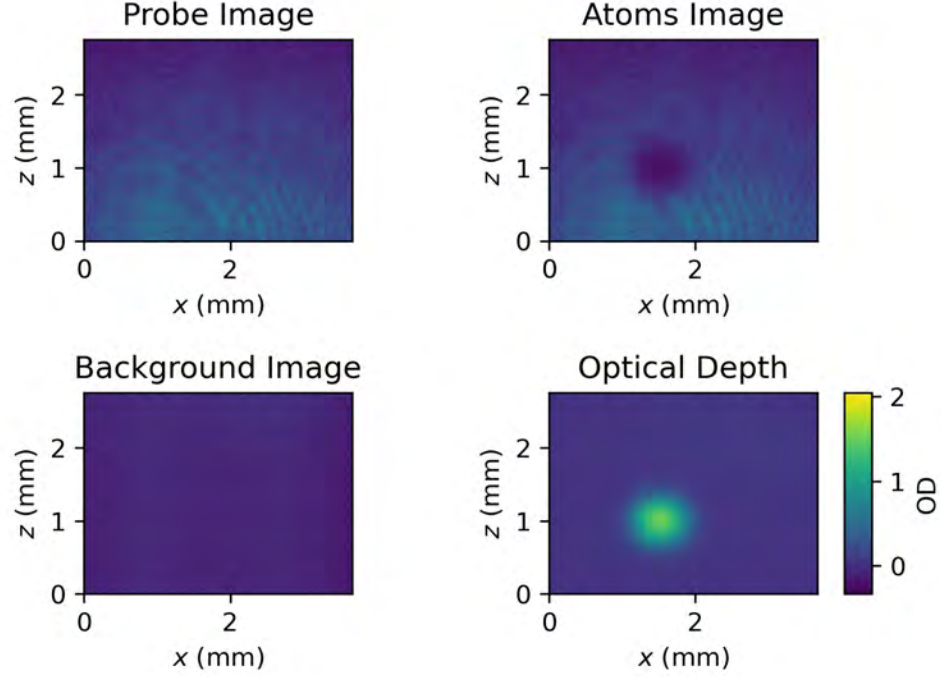


Figure 4.1: An example of the three images we take when doing absorption imaging, and the optical depth they combine to create. The probe image is in the upper left which is taken once the atoms have left the frame. The atoms image is in the upper right, and you can see the shadow of the atomic cloud. The background image is in the lower left. Finally the optical depth, the combination of these three images, is in the lower right.

a beam propagating a distance dz through a cloud of atoms with density n . If we take $I \ll I_s$ this drop in intensity becomes [1]:

$$dI = -n(x, y, z)\sigma_0 \frac{I}{1 + (2\delta/\Gamma)^2} dZ$$

where σ_0 is the on resonance cross section, and assuming a two level system $\sigma_0 =$

$3\lambda^2/2\pi$. If we take $\delta = 0$ and integrate over z we are left with

$$-\ln\left(\frac{I_f}{I_i}\right) = -\ln\left(\frac{I_{\text{atoms}} - I_{\text{bg}}}{I_{\text{probe}} - I_{\text{bg}}}\right) = n_{2D}(x, y)\sigma_0,$$

Where n_{2D} is the density of atoms integrated along z . We then define the Optical Depth (OD) as $n_{2D}\sigma_0$. Adjustments to this analysis have to be considered when probing with higher powers, with detuned beams, and when dealing with something other than a simple two-level system. For the bosonic isotopes of strontium the extension beyond a two-level system does not change much, but the fermion requires a more complicated analysis of the absorption cross section, described in detail in Ref. [56]. This analysis shows that the highest cross section, and best absorption, occurs with a circularly polarized probe beam for ^{87}Sr . Once the OD is calculated getting the atom number is a simple matter of integrating the OD over the extent of the cloud in x and y .

If we assume the atoms in the trap are in thermal equilibrium, then the velocity distribution in one dimension is a Maxwell-Boltzmann distribution:

$$f(v_x) = \left(\frac{m}{2\pi k_b T}\right)^{1/2} e^{\frac{-mv_x^2}{2k_b T}},$$

where k_b is the Boltzmann constant. As the cloud of atoms expands in time of flight the spacial distribution will be a convolution of this distribution with the initial spacial distribution of atoms in the trap. Thus we generally fit the images to a Gaussian function, which becomes a better assumption as the time of flight

increases. Ideally we would fit to a 2D Gaussian function, but this is generally slow, and we have found fitting to 1D slices in both dimensions of the image to be efficient and accurate. We do an iterative fit, where the slice of the image we fit to along one axis is determined by the center found from the previous fit to the orthogonal axis. This ensures the fits find the true center of the cloud in the 2D image.

Once we are able to take images in time of flight we can calibrate the magnification of the imaging system by dropping the MOT, and watching it fall. By measuring the acceleration of the MOT downwards, and matching it to gravity we can convert between pixels on the camera and meters. Currently, the horizontal imaging system on the main apparatus has a conversion of 3.4×10^{-6} m/pixel (magnification of 1.4), and the imaging system on the grating MOT apparatus has a conversion of 19.4×10^{-6} m/pixel (magnification of 0.25). Looking at the MOT after a time of flight also gives us a measure of the temperature. After a time t the width of the cloud becomes

$$w^2 = w_0^2 + \frac{k_b T}{m} t^2, \quad (4.1)$$

where w_0 is the initial width of the cloud. Thus fitting the MOT expansion allows us to extract the temperature of the cloud in both axes imaged by the camera.

4.2 Blue MOT

The blue MOT of the main apparatus is relatively simple. We operate at a field gradient around 55 G/cm, and with about 9 mW of power in the MOT beams, giving

a peak I/I_{sat} of ~ 0.4 . After loading the trap for 2 to 4 s we trap $\sim 3 \times 10^7$ atoms of ^{88}Sr , and for the other isotopes we trap atoms approximately proportional to the atomic abundances. For ^{87}Sr , the 707 nm repump EOM is modulated at a variety of frequencies to account for the hyperfine structure. In the main apparatus we modulate at 571 MHz, 1240 MHz, and 480 MHz. These frequencies are empirically found informed by the values from the old laboratory [5]. Due to the magnification of our imaging system and relatively high temperature of the blue MOT, in time of flight the atomic cloud quickly expands past the field of view. This makes it difficult to take a good measurement of the temperature of the atoms in the blue MOT, but we expect that it is somewhere around a few mK.

After loading the blue MOT we ramp the power in the MOT beams down, while ramping the field gradient up. We call this section of the blue MOT experimental sequence the “Doppler cooling” stage. Fig. 4.3 shows the laser power and field ramp during this stage. Lowering the power should allow us to reach a temperature that is closer to the Doppler limit, which is only achievable in the limit of low saturation, while ramping the field up compresses the MOT, counteracting the expansion of the MOT that occurs when the power is lowered. Lowering the temperature facilitates better transfer between the red and blue MOT. In practice, we currently get the most efficient transfer when the blue power is dropped to zero, and the field barely increases, although the transfer is not very sensitive to the field values, so they can change depending on the specific optimization.

4.3 Red MOT

Finding the red MOT is quite a bit more complicated than the blue MOT. Firstly, since the red MOT requires first loading the blue MOT, and the red MOT is not visible by eye, we cannot continuously load and optimize the MOT at the same time, and we have to use destructive imaging to analyze the red MOT. Second, there must be significant spatial overlap between the red and blue MOT in order for there to be sufficient transfer. Furthermore, the lower field gradient of the red MOT causes any stray field to considerably move the field zero. This means we can move the red MOT around in space easily with the shim fields, but if the fields are not correct the red MOT field zero can easily be in a completely different location than the blue MOT, and transfer will be impossible.

We have found a relatively consistent way to find the ^{88}Sr red MOT with our main apparatus. First, there is a visible dip in fluorescence and a change in the shape of the blue MOT from the red MOT beams when they are on resonance. Since this is visible even while continuously imaging the fluorescence of the blue MOT, it allows us to easily confirm the laser is close to resonance. Once we are close, we leave the red light on while looking at absorption images of the blue MOT in time of flight with the field gradient set to an appropriate red MOT value (between 3 and 5 G/cm). After a few ms only a diffuse background gas from the blue MOT is visible. Atoms in the background gas see a force towards the field zero, and are excited to the 3P_1 level wherever the light is resonant with the Zeeman shifted transition. Since the linewidth of the red MOT transition is so narrow this condition

is only satisfied for a thin shell of atoms where the field is appropriate, and these shells appear in the images as ellipses. We can look at the center of these shells and how they move with changing frequency to adjust the shim fields so that the center of the shells is overlapped with the blue MOT. This also tells us where our eventual red MOT will be, so we can tune other parameters until the MOT appears at that spot. Additionally, the radii of these ellipses tell us the approximate detuning of the red light. Fig. 4.2 show example images of these shells.

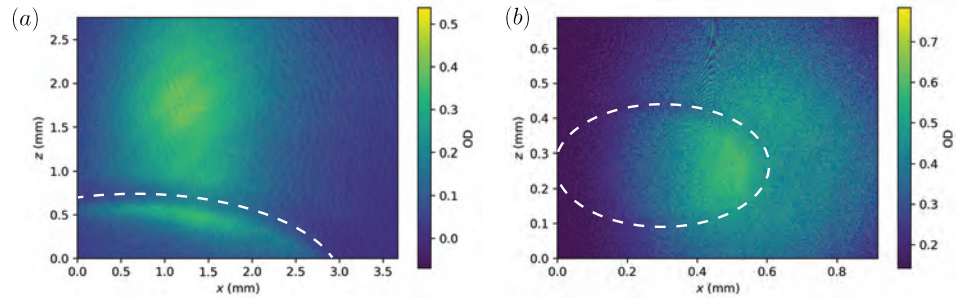


Figure 4.2: Absorption image of the 689 nm light interacting with the blue MOT background. The resonant shells are marked with a dashed line. In (a) we can see that the center of the resonant shell is well below the blue MOT, so we have to use the shim fields to bring it up. Image (b) shows the shell better aligned with the blue MOT. From here we are likely to get a red MOT with just a few adjustments of other parameters.

We start loading the red MOT with an artificially broadened laser, about 2 MHz wide, and slowly ramp it down to a single frequency. Broadening the laser enables us to initially capture more velocity classes from the blue MOT and greatly increases the final atom number. While we ramp down the broadening we also ramp down the power in the red MOT beams to further cool the MOT using the RF power to the AOM before the red fiber splitter. Fig. 4.3 shows example parameters for the

blue to red MOT transfer. We have captured $\sim 1 \times 10^7$ ^{88}Sr atoms in the red MOT, at a temperature of $\sim 5 \mu\text{K}$. We have also achieved a red MOT of all other stable isotopes, with atom numbers consistent with the relative abundances. For ^{87}Sr the stirring laser, mentioned in Sec. 2.1.3, is broadened in the same way, and in phase with the cooling laser, and the power is ramped down similarly as well, this keeps the stirring continuously working while the broadening is narrowed.

4.4 Dipole Traps

The main difficulty with transitioning from the red MOT to the dipole trap is aligning the two in space. We want the *in situ* red MOT to sit right where the focus of the dipole beam or beams are. To get the initial alignment we first set up a temporary imaging system, along the axis of the dipole trap, but on the other side of the chamber from the input dipole beam. Then we add in a flipper mirror so we are able to switch between putting the dipole trap and a probe beam through this imaging system. Then we take *in situ* images of the red MOT which allows us to focus the imaging system, and note the location of the red MOT in the images. We then swap in the dipole beam for the probe beam, and use the same imaging system to align the image of the dipole beam up with the previous image of the MOT to give approximate transverse alignment. The trap can be aligned axially by adjusting the final lens of the dipole trap to focus it on the camera. Since we previously focused the imaging system with the red MOT, this process aligns the focus of the dipole trap onto the atoms. Throughout this process we lower the

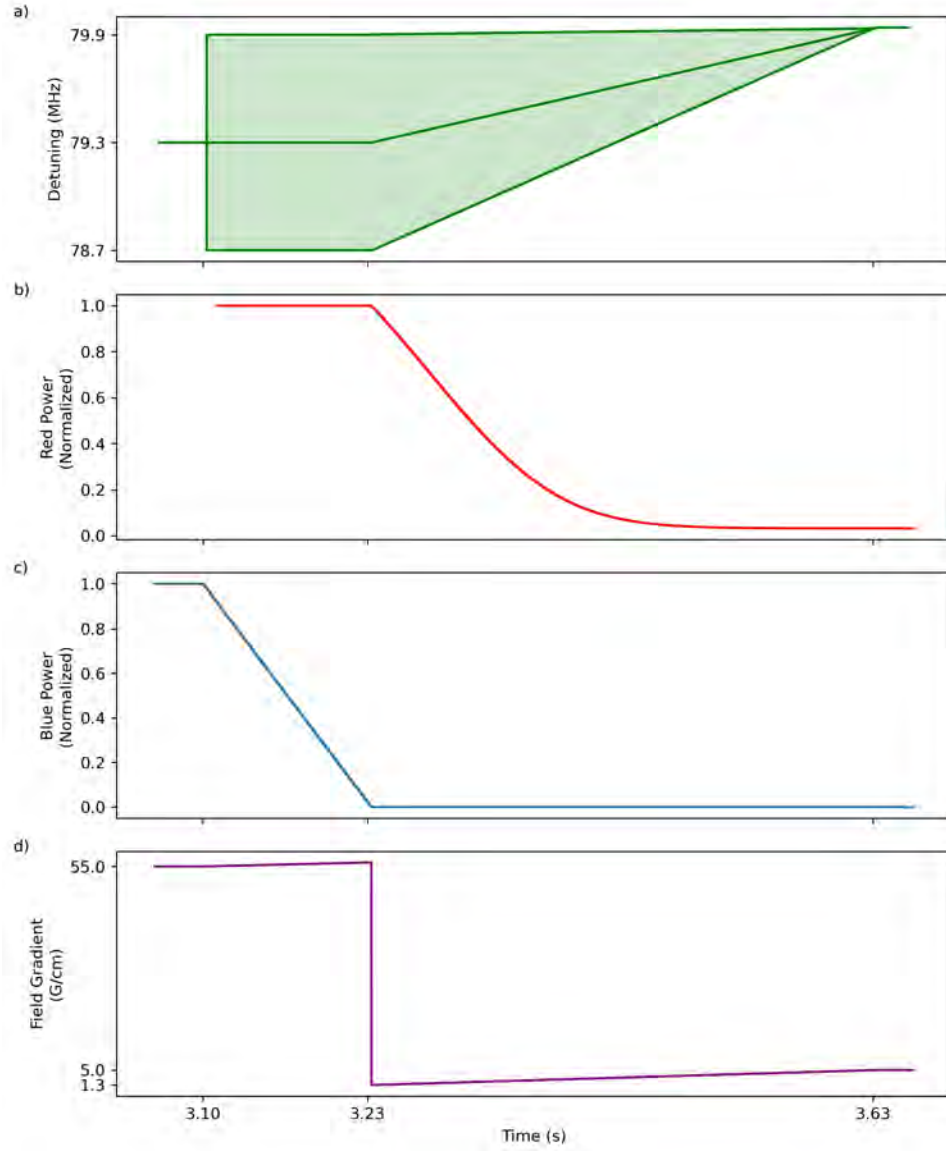


Figure 4.3: Example red MOT sequences for the MOT in the main apparatus. a) The broadening of the laser detuning starts out large and decreases to zero over time. b) The power in the red MOT beam decreases as well to cool the atoms in the MOT to the final temperature. We ramp this power with the AOM RF power linearly which leads to the exponential shape of the optical power ramp. c) The blue power decreases during the Doppler cooling stage using the power lock, and d) the field gradient increases slightly. After the Doppler cooling stage the field gradient drops to the initial red MOT value where it is slowly ramped up.

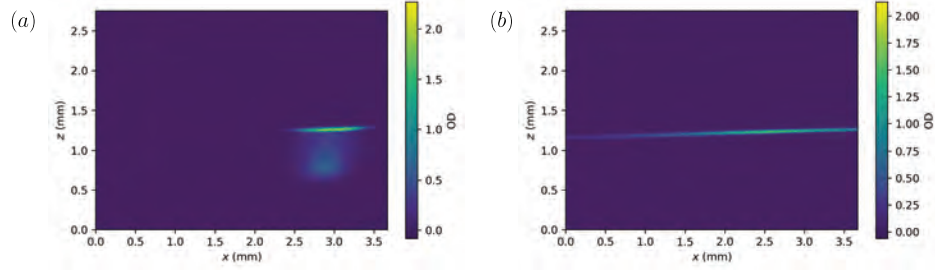


Figure 4.4: (a) One of the first images we took of our dipole trap 10 ms after the red MOT was dropped. You can see the red MOT atoms that were not captured by the dipole trap continuing to fall. (b) The dipole trap after some optimization and after holding the dipole trap for 5 s. You can see the weak trapping in the axial direction, and the slight angle to the trap that still needs to be leveled.

power in the dipole beam significantly, and add a neutral density filter to avoid any damage to the camera. While this process gives us rough alignment, chromatic shifts in the imaging system between the probe wavelength (461 nm) and the dipole trap wavelength (1064 nm in this example) can give a significant offset from perfect alignment. We make up for this final misalignment by randomly sampling a large space for the final red MOT shim coil values. This moves the red MOT around in space with a very good chance of overlapping the dipole beam at some point. You can see the result of this alignment in Fig. 4.4. After this process the red MOT shim fields might not be at optimal values for the blue to red MOT transfer, but once the initial trap is found it is simple to slowly walk the dipole beam alignment and shim values back to optimal values.

For many applications a single beam dipole trap is not sufficient, and two (or more) beams, intersecting at a point, must be used. The process for aligning this second, cross dipole beam is the same as the first, we do this alignment without

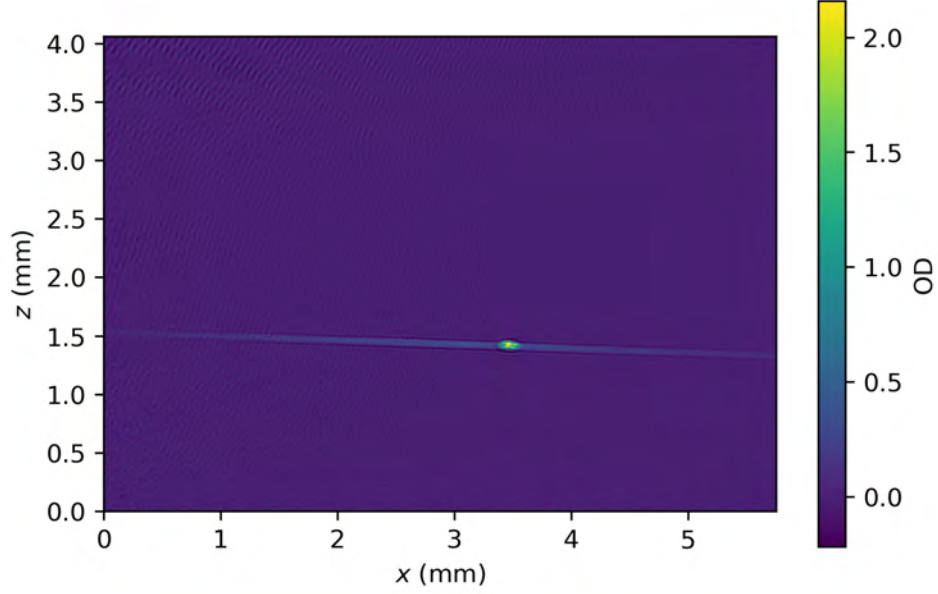


Figure 4.5: An image of the atoms in the cross dipole trap with two dipole trap beams. Note that the magnification of our imaging system changed between this image and those in Fig 4.4.

operating the main dipole trap. Once we move the red MOT so that some atoms were trapped by the cross dipole beam we have two different sets shim values for the two dipole beams. All that is left is to walk the main dipole beam's shim values to the values for the cross beam so that both beams and the red MOT all overlap. Our cross beam comes in on the same axis as our permanent horizontal imaging system, so a permanent dichroic is used to combine the probe and dipole beams. However, since the main probe beam and cross dipole beam counter-propagate through the chamber, a temporary imaging system is still required on the opposite side of the chamber from our permanent imaging system.

The traps discussed here are 1064 nm dipole traps, but for most of our applications we want magic wavelength 813 nm attractive traps. This initial alignment

should be sufficient to easily swap out the light for the magic wavelength and still see the trap, but there will be some chromatic shifts which will impact the quality of the trap. Thus, we did not invest time in fully optimizing the 1064 nm dipole traps while we did not have access to 813 nm light.

4.5 Blue Grating MOT

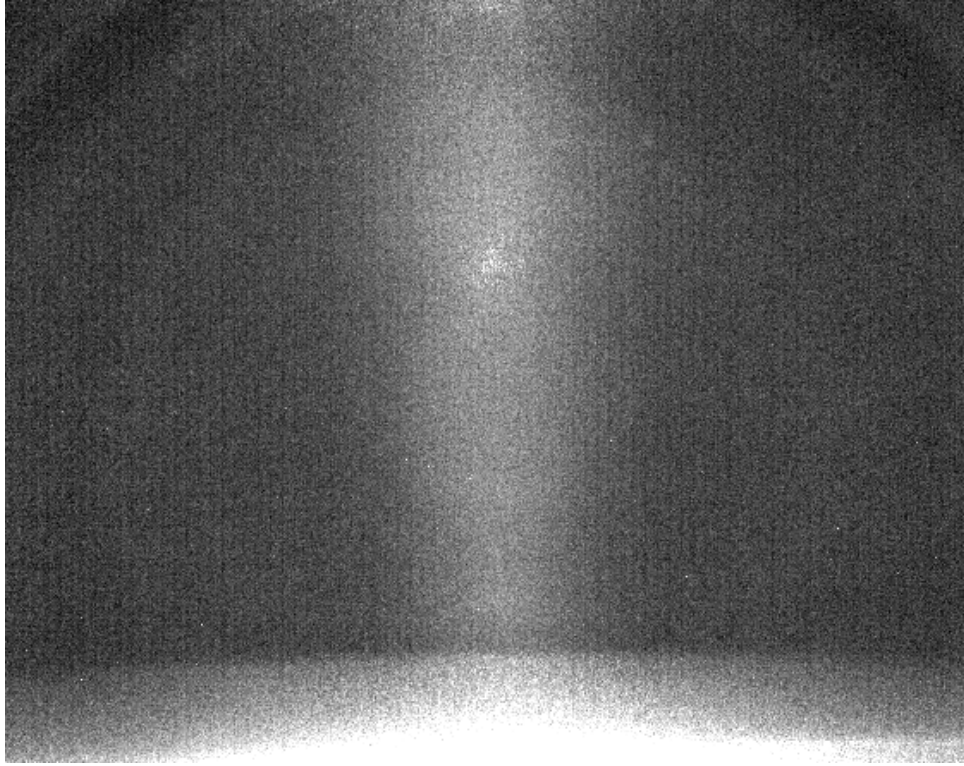


Figure 4.6: An image of the atomic beam out of the dispenser source with the magnetic fields off and the MOT beams on resonance.

Just like the blue MOT of the main apparatus, the blue grating MOT is relatively simple to find. The atomic fluorescence of the atomic beam is visible when the magnetic field is off, and the MOT beam is on resonance, as seen in Fig.

4.6. This allows for the resonance to be found independent of alignment and the magnetic field. Once the field is on, the easiest way to align the MOT beam is to look at the 0th order diffraction of the grating which should retro-reflect back onto the main beam. In practice however, just like most MOTs, perfect retro-reflection is not optimal. The physical z position of the coils or magnets can also be adjusted to optimize the MOT.

While the MOT is loading we take fluorescence images approximately every 200 ms. This allows us to see the MOT loading rate. Unfortunately, due to some noise in the blue MOT laser lock, the fluorescence sometimes instantaneously drops, but in general there are enough usable images. Fig. 5.6 shows examples of these loading curves. From the fits to these curves we can calculate the final relative atom number, unfortunately, it is difficult to get an accurate absolute atom number using fluorescence imaging. Additionally, in the first iteration of the blue grating MOT, absorption imaging was only possible when the atom number was optimal, which restricted the data that could be taken with this method, and added significant noise to the data we could take. To combat this issue, we normalized the fluorescence imaging to the absorption atom number using a constant averaged normalization factor. This allowed us to get a stable and accurate atom number from the loading curves. Since the failure of the atomic shutter prohibits us from shutting off the atom beam and looking at the decay of the MOT, we instead use the loading curves to extract the lifetime of the MOT. Sec. 5.2.4 discusses more details on these loading curves and their interpretation. Since the fluorescence in the loading images is from the blue MOT beam, in order to create Fig. 5.7, we had to account for the

dependence of fluorescence on detuning. Luckily this is simple using the equation for the scattering rate (Eq. 2.1).

Since the imaging system on the grating MOT apparatus demagnifies, compared to the magnification of the main apparatus imaging system, we get a larger field of view. This allows us to take images with varying time of flight, and measure the temperature of the MOT. This process is described in Sec. 5.2.4.

After the publication of Ref. [58] (Sec. 5.2), which did not include any results from the fermion ^{87}Sr , we found the fermionic blue grating MOT. By adjusting the magnets to account for the differential offset in position between the fermion and the bosons as discussed in Sec. 2.3 we were able to trap the fermion. For the fermionic grating MOT, the 707 nm repump is modulated at 1315 MHz and 499 MHz, these values are informed by Ref. [69]. Additionally, in the second iteration of the grating MOT apparatus, the atom number in the blue MOT was improved by an order of magnitude, up to 3×10^7 . This was likely due to using the same blue laser as the main apparatus, which has more available power, for the MOT beam, and an improvement in the vacuum pressure down to approximately 1×10^{-10} Torr. This increase in the atom number makes absorption imaging a much more consistent tool.

To transfer atoms from the blue grating MOT to the red grating MOT we operate the same Doppler cooling stage as on the main apparatus. However, since we are currently operating the blue grating MOT coils at the maximum current of the power supply, we cannot ramp the field gradient up during this stage. In the grating MOT the length of this stage that gives an optimal transfer is much shorter than in the main apparatus, as seen in figure 4.8, this could be because we cannot

compress the MOT with a field gradient.

4.6 Red Grating MOT

The red grating MOT has difficulties that are not present in the conventional red MOT. The capture region is smaller and does not overlap completely with the blue capture region due to the difference in diffraction angle. The high transverse temperature of the blue grating MOT means that atoms can quickly leave this capture region. Additionally, before the apparatus was refined, the slow rate of field switching further exasperated this issue as it left more time for the atoms to expand before the red MOT field value was reached.

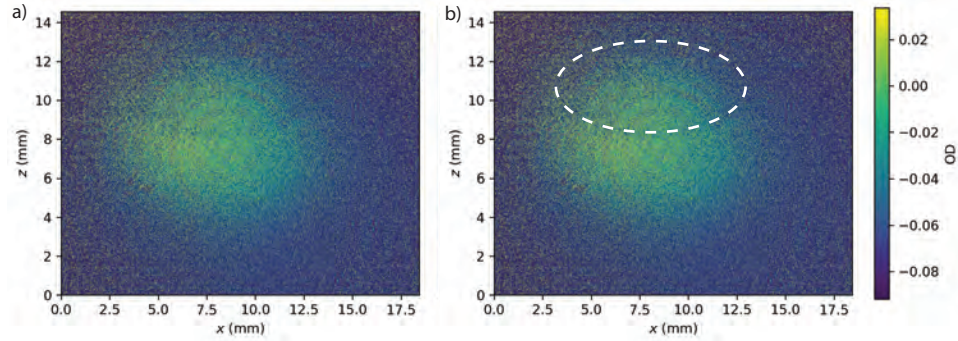


Figure 4.7: a) An image of the resonant shell from the red MOT beams in the background of the blue grating MOT. b) The same image with the shell marked with a dashed white line.

The comparatively low atom number in the blue grating MOT also adds difficulty by decreasing our signal to noise ratio across the board. For example, the signal of the resonant shells discussed in Sec. 4.3, is greatly reduced, especially at low field gradient, and they are less useful as a diagnostic tool. An example image

of a resonant ring in the blue grating MOT is shown in Fig. 4.7. Despite this, the rings were useful in diagnosing the issues we had with the field switching. From the relative insensitivity of the radii of the rings to the red laser detuning, we could determine that our field gradient was still high even 2 ms after commanding a low field. Additionally, we could tell by how the rings shifted in the $+z$ direction over time that we had an eddy current from the grating mount.

Once the issues with field switching were resolved, as discussed in Sec. 3.1.4, the main method we used to search for the red MOT was a large random sampling of the appropriate parameter space to search for any initial signal. In this way, we found a very faint signal of diffuse atoms at 10 ms time of flight (longer than blue MOT atoms remain visible), and from here a blind adjustment of alignment got us our first red Grating MOT.

Example parameters for the transfer between the blue and red grating MOTs are shown in Fig. 4.8. Here you can see the short Doppler time where the blue MOT power decreases to zero, the frequency broadening and compression, the decrease in the red MOT power to compensate for this compression, and the field switching. The broadening in the red grating MOT is done using a sawtooth function to facilitate SWAP cooling. In addition to perform SWAP we increase the broadening, which pushes the top edge of the broadening above resonance.

In the red grating MOT we can capture around 3×10^6 atoms, achieve a temperature of around $3 \mu\text{K}$, with a trap lifetime of 0.8 s. The measurements of temperature and lifetime are shown in Fig. 4.9 and Fig. 4.10 respectively.

We also investigate the effect of SWAP, discussed in Sec. 2.4, on the grat-

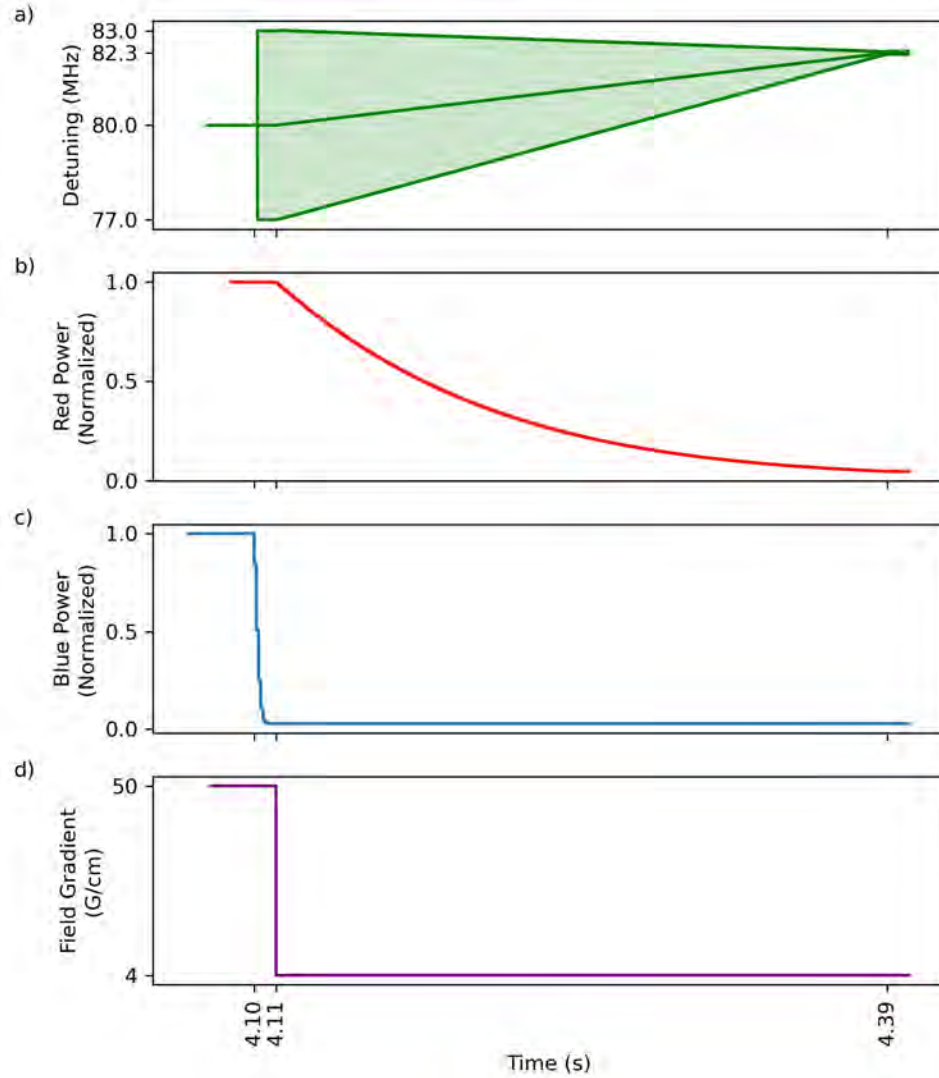


Figure 4.8: Example red MOT sequences for the grating MOT. a) In the grating MOT the broadening is much larger at first to account for the higher temperature of the blue grating MOT, and to facilitate SWAP. In addition, for SWAP the frequency of the upper edge of this broadening decreases slightly over time, and starts above the resonance, and the final broadening is left at 100 kHz rather than zero. b) The corresponding decrease in red MOT beam power is shown. c) The blue power decreases during the Doppler cooling stage (linearly in AOM RF power), but in the grating MOT the length of this stage is much shorter. d) The field gradient switches from the blue MOT value to the red MOT value, unfortunately these are the maximum operating values for the respective coils, so we cannot ramp the field in either the Doppler stage or broad red MOT stage like we do on the main apparatus.

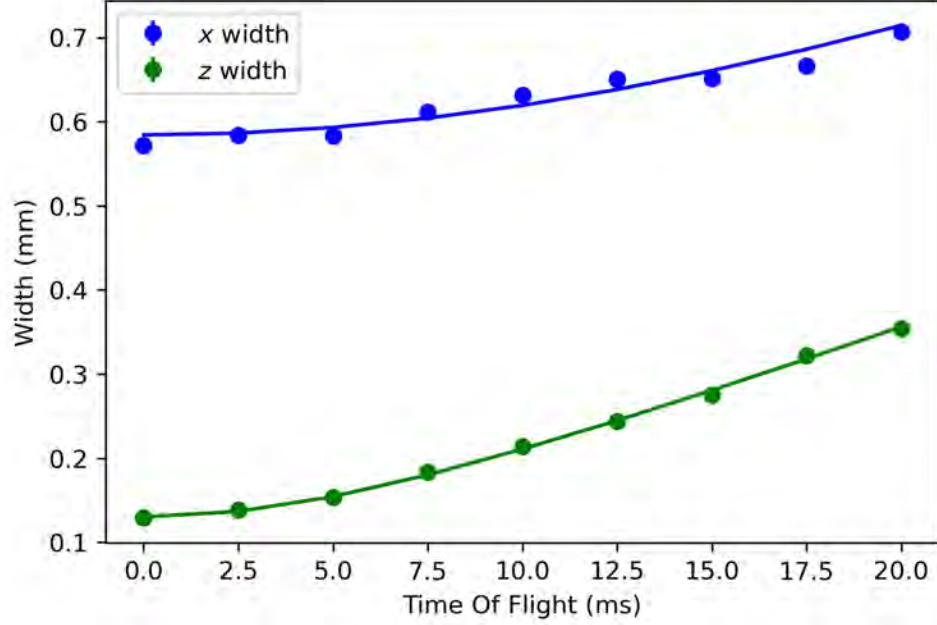


Figure 4.9: A measurement of the red grating MOT temperature. We measure the temperature by looking at how the width of the cloud increases in time of flight. The fits to this expansion, using Eq. 4.1, are shown as well. These fits yield temperatures of $4.5(6) \mu\text{K}$ in the transverse (x) direction, and $2.9(1) \mu\text{K}$ in the axial (z) direction.

ing MOT by switching between sawtooth and triangle sweeps, as well as sawtooth sweeps with the wrong slope (anti-SWAP). Fig. 4.11 shows how the atom number varies between these three different methods of modulation, and you can see the enhancement in atom number that SWAP provides. SWAP at 20 kHz modulation almost doubles our atom number when compared to the optimal triangle wave modulation at 20 kHz. We can match the sweep rate of the triangle and sawtooth waves by halving the frequency of the triangle sweep to 10 kHz, but this simply drops the atom number from the triangle modulation further. Lastly, modulating with a sawtooth wave with the wrong slope pushes atoms out of the trap, and drastically

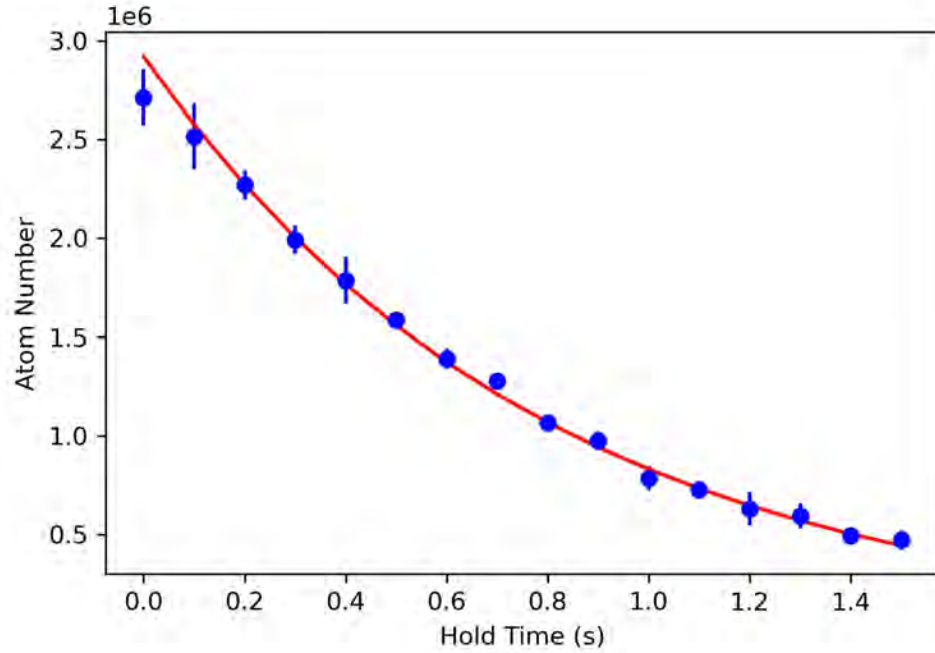


Figure 4.10: A measurement of the red grating MOT lifetime. A fit to the exponential decay of atoms in the trap as a function of hold time is shown in red. This yields a lifetime of 0.8 s.

drops the atom number. In addition Fig 4.12 shows how the atom number varies as a function of the sweep frequency.

In conclusion, on the main apparatus we got our new strontium source up and running, and transferred atoms between the progressive cooling stages of strontium from the source, to the blue MOT, to the red MOT, to dipole traps. We have improved the atom number in the red MOT compared to the old apparatus. The next step for the main apparatus involves implementing magic wavelength dipole traps and preparing for future experiments discussed in Ch. 6. On our grating MOT apparatus we implemented the first alkaline earth grating MOT, performed transfer to the red grating MOT, and showed enhancement of the transfer efficiency with

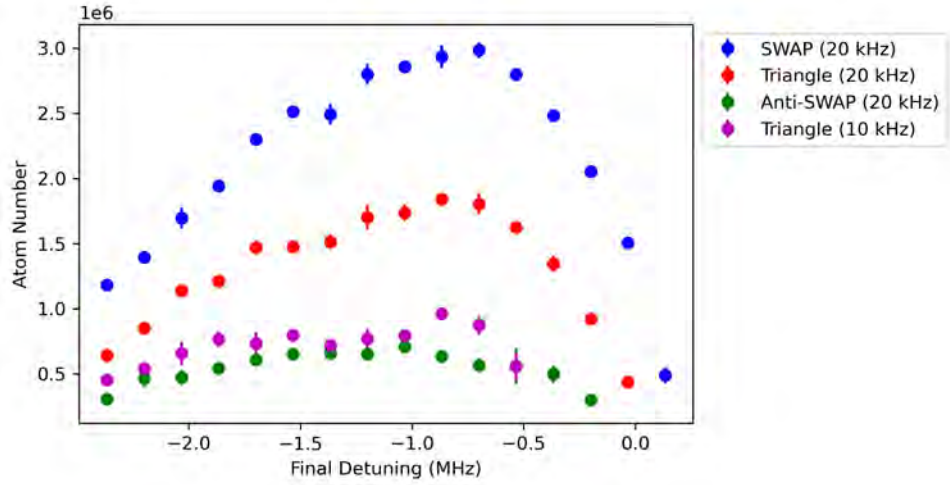


Figure 4.11: Atom number vs detuning for SWAP vs a triangle wave sweep. The detuning shown is just an estimation of the final detuning of the laser beam after the broadening is decreased down to 100 kHz. The data for our optimal SWAP MOT is shown in blue with a sawtooth wave at 20 kHz. The standard swept MOT with a triangle wave is shown in red, and a sawtooth wave MOT where the sawtooth has the wrong slope (anti-SWAP) is shown in green, both with a 20 kHz sweep frequency. We also show the triangle wave MOT with a 10 kHz sweep frequency, shown in purple, to match the sweep rate of the SWAP MOT (note that these data were taken a different day so they were normalized to 20 kHz SWAP MOT).

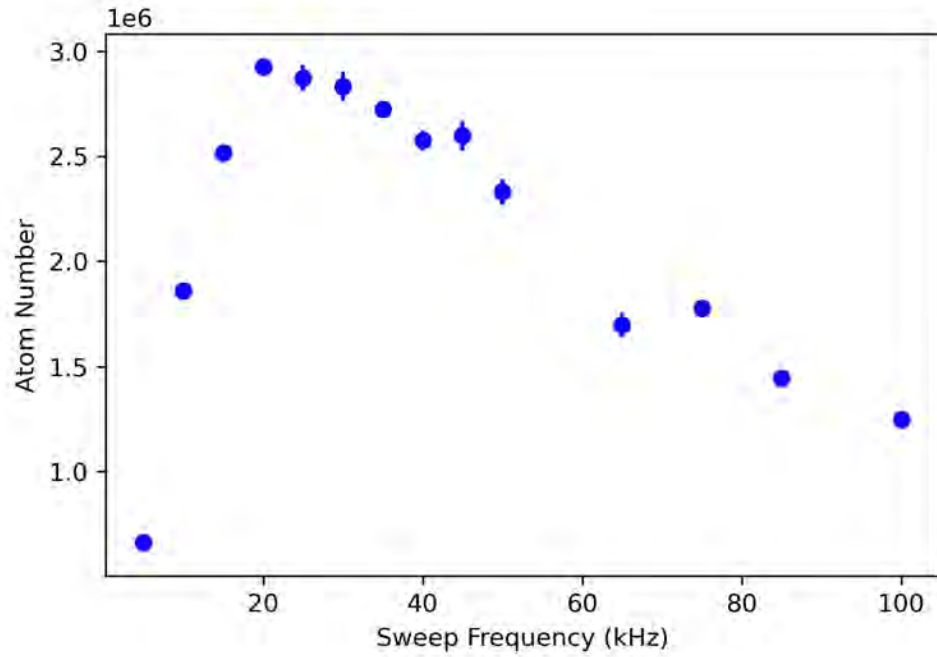


Figure 4.12: The atom number we can achieve through SWAP at a variety of different sweep frequencies. Our atom number is optimal when the sweep frequency is 20 kHz.

SWAP in the grating geometry. The immediate next step for the grating apparatus is to find a red MOT of ^{87}Sr .

Chapter 5: Publications

5.1 Publication: Isotope shift spectroscopy of the $^1S_0 \rightarrow ^3P_1$ and $^1S_0 \rightarrow ^3P_0$ transitions in strontium

This work was completed in the old strontium laboratory, published as Ref. [24] and led by Neal Piseni. Isotope shifts were a straightforward avenue to investigate given the laboratory's capability to operate with all the stable isotopes in strontium, and the lack of precision spectroscopy of the narrow transitions in the less used isotopes. In addition with four stable isotopes, we had enough data points to investigate the linearity of the results, which became more interesting given recent proposals to use this linearity as a way to probe new physics. The precision spectroscopy was further motivated by the recent purchase of an optical frequency comb by the JQI. Halfway through the data collection, Neal moved on to other projects and I took over the work in the laboratory for the second half of the experiment, specifically the spectroscopy and systematics of the clock transitions. Much of the text was written by our postdoc Hiro Miyake, with significant input from the rest of the group. This work found a significant deviation from linearity in the king plot due to the hyperfine structure of ^{87}Sr , and work since has provided more possible

explanations for this effect [70–72].

5.1.1 Abstract

Isotope shift spectroscopy with narrow optical transitions provides a benchmark for atomic structure calculations and has also been proposed as a way to constrain theories predicting physics beyond the Standard Model. Here, we have measured frequency shifts of the $^1S_0 \rightarrow ^3P_1$ and $^1S_0 \rightarrow ^3P_0$ transitions between all stable isotopes of strontium relative to ^{88}Sr . This includes the first reported measurements of the $^1S_0 \rightarrow ^3P_0$ isotope shift of ^{88}Sr - ^{86}Sr and ^{88}Sr - ^{84}Sr . Using the isotope shift measurements of the two transitions a King plot analysis is performed, revealing a non-linearity in the measured values.

5.1.2 Introduction

Isotope shifts of atomic transition frequencies arise due to the difference in neutron numbers for different isotopes with the same atomic number. For a given element, these shifts can be systematically analyzed using a King plot, which elucidates the contributions of the field and mass shifts [25]. The King plot is typically expected to be linear, and the experimentally determined value of the slope provides a good benchmark for theoretical predictions [27]. Any deviations from linearity as was observed in Sm [73] and Ba [74], or between predicted and experimentally measured values of the slope as was observed in Ca^+ [75], are important for refining atomic structure calculations [76]. Furthermore, recent theoretical proposals have

suggested that linearity in King plots could be used to put constraints on higher-order effects on isotope shifts or on physics beyond the Standard Model [26, 77]. Strontium has many favorable properties for studying isotope shifts, including an abundance of stable isotopes and very narrow optical transitions [40]. In addition, prior theoretical work has proposed the measurement of strontium isotope shifts as a promising probe of new physics [26, 77].

Strontium has four stable isotopes: three bosons (^{88}Sr , ^{86}Sr , and ^{84}Sr), and one fermion (^{87}Sr). Mixing between the singlet and triplet fine structure manifolds leads to narrow-linewidth optical transitions, and these transitions have found use in both strontium and other alkaline-earth-(like) atom experiments [78, 79]. In particular for strontium, the $^1S_0 \rightarrow ^3P_1$ intercombination-line transition at 689 nm (linewidth $\Gamma/2\pi = 7.4$ kHz) is used during laser cooling to operate a narrow-line magneto-optical trap (MOT) [40, 80], and the even narrower $^1S_0 \rightarrow ^3P_0$ clock transition at 698 nm ($\Gamma/2\pi \sim$ mHz) is the foundation for state-of-the-art optical clocks operating at a precision at the 10^{-18} level [11, 12, 81]. The clock transition is strictly forbidden by angular momentum considerations, but becomes weakly allowed via hyperfine mixing in ^{87}Sr or by application of an external field for the bosonic isotopes [9].

While the $^1S_0 \rightarrow ^3P_0$ clock transition has been extensively studied in ^{87}Sr and ^{88}Sr [9, 11, 12, 81–85], there have been no previous measurements of the transition in either ^{86}Sr or ^{84}Sr as far as we know [86]. Here we report the first isotope shift spectroscopy measurements of the clock transition for both ^{84}Sr and ^{86}Sr relative to the most abundant isotope ^{88}Sr . Furthermore, we measure all isotope shifts of the $^1S_0 \rightarrow ^3P_1$ intercombination-line transition relative to ^{88}Sr , permitting the first King

plot analysis of strontium for these two transitions. Given the very narrow linewidths involved, extensions of this work could place stringent experimental constraints on the King linearity, ruling out candidate theories for physics beyond the Standard Model or benchmarking state-of-the-art atomic structure calculations.

5.1.3 Experimental procedure

All of the isotope shift spectroscopy was performed using laser-cooled strontium atoms at temperatures of a few μK , held in an optical dipole trap (ODT). After applying the spectroscopy light, we monitored atom loss by performing absorption imaging.

	88	87			86	84
Systematic Shift (kHz)		$F' = 7/2$	$F' = 9/2$	$F' = 11/2$		
Density	1.7 ± 1.7	-34.1 ± 14.5	-51.9 ± 26.8	-43.3 ± 15.6	5.1 ± 3.4	-1.4 ± 4.3
Recoil	$4.8 \pm (< 0.1)$	$4.8 \pm (< 0.1)$	$4.8 \pm (< 0.1)$	$4.8 \pm (< 0.1)$	$4.9 \pm (< 0.1)$	$5.0 \pm (< 0.1)$
Total	6.5 ± 1.7	-29.3 ± 14.5	-47.1 ± 26.8	-38.5 ± 15.6	10.0 ± 3.4	3.6 ± 4.3

Table 5.1: Measured systematic frequency shifts and uncertainties for the $^1S_0 \rightarrow ^3P_1$ transition. Uncertainties indicate one standard deviation.

The laser lights used for spectroscopy of both the $^1S_0 \rightarrow ^3P_1$ and $^1S_0 \rightarrow ^3P_0$ lines were generated using two home-built external-cavity diode lasers based on the design in Ref. [64]. The frequency of the 689-nm laser was stabilized via an optical phase-locked loop [87] to the master laser of the 689-nm narrow-line MOT system. The master 689-nm laser was locked using the Pound-Drever-Hall (PDH) method [88, 89] to a cavity constructed from ultra-low expansion (ULE) glass and housed in a temperature-stabilized vacuum chamber. To stabilize the frequency of the 698-nm laser, we passed a few percent of the light through a wide bandwidth electro-optic modulator and locked the first phase-modulated sideband

via PDH to a second, independent ULE cavity [90]. This cavity was housed in a separate temperature-stabilized, acoustically-isolated vacuum chamber. These locking schemes for both lasers allowed us the flexibility to shift the frequency of either laser to span the isotope shifts of its respective transition. For both the 689-nm and 698-nm lasers, the light was referenced to a frequency comb (Menlo System FC1500-250-ULN [91]) to account for long-term drift and provide a frequency reference. Fine frequency control of each laser beam was achieved by adjusting the drive of an acousto-optic modulator, which was also used to stabilize the intensity of the spectroscopy pulse. The spectroscopy laser linewidth was characterized by locking independent 689-nm lasers to each ULE cavity. A heterodyne beatnote at 689 nm between the two separate lasers was measured to be approximately 200 Hz wide, which bounds the expected spectral performance of both systems.

The remainder of the apparatus used for the spectroscopy has been described in detail previously [43]. Laser cooling of all isotopes proceeds according to well-established techniques [40], with a MOT first operating on the broad $^1S_0 \rightarrow ^1P_1$ transition at 461 nm, followed by a narrow-line MOT operating on the 689-nm intercombination-line transition $^1S_0 \rightarrow ^3P_1$. For all isotopes, temperatures in the narrow-line MOT are typically a few μK , low enough to efficiently transfer the atoms into a single beam, far-detuned ODT at 1064 nm. Typical temperatures in the ODT are $\{2.9 \mu\text{K}, 2.2 \mu\text{K}, 1.1 \mu\text{K}, 2.7 \mu\text{K}\}$ and typical atom numbers are $\{1, 0.1, 0.5, 0.2\} \times 10^6$ for $\{^{88}\text{Sr}, ^{87}\text{Sr}, ^{86}\text{Sr}, ^{84}\text{Sr}\}$ respectively, with trap frequencies $\{\omega_x, \omega_y, \omega_z\}/(2\pi) = \{50, 4, 495\}$ Hz in the horizontal, axial, and vertical directions respectively. For the bosons, the variation in atom number is mostly due to the

difference in the natural abundance of each isotope, whereas for the fermionic isotope the atom number is also limited by the additional complexity of the narrow-line MOT [44].

5.1.4 Measurement of the $^1S_0 \rightarrow ^3P_1$ isotope shifts

After loading the atoms into the ODT, the magnetic field was set to 0.05 mT (0.5 G) to resolve the $^3P_1(m = 0)$ state for the even (bosonic) isotopes. For the odd (fermionic) isotope, which has hyperfine structure, the magnetic field was set to zero, meaning that the Zeeman splitting was not detectable within the lineshape. The strength of the magnetic field was calibrated by addressing the $^1S_0(m = 0) \rightarrow ^3P_1(m' = 1)$ transition of ^{88}Sr , for which the Zeeman shift is known [80]. To eliminate the effect of AC Stark shifts from the ODT we implemented a stroboscopic procedure where, with a typical period of 500 μs , the ODT was turned on and off with a duty cycle of 50% (duration of 250 μs), and applied the 689-nm probe laser when the ODT was off, similar to the procedure used in Refs. [92, 93]. The spectroscopy light was used to induce atom loss from the trap through light scattering and subsequent recoil, which is primarily an incoherent process where there is absence of coherence between the pulses of the stroboscopic method with the switching of the ODT. The 689-nm spectroscopy beam was aligned at an angle of approximately 45° with respect to the ODT, both in the horizontal plane. The spectroscopy beam was collimated with a $1/e^2$ beam waist of 1.25 mm in the horizontal direction and 1.71 mm in the vertical direction at the position of the atoms. The polarization of the

spectroscopy beam was set to be linear along the direction of the magnetic field. The total illumination duration used for spectroscopy was set to between 1 ms and 15 ms, corresponding to multiple stroboscopic pulses, and the peak optical intensity was at most 0.1 mW/cm^2 ($I_{\text{sat}} = 3 \text{ }\mu\text{W/cm}^2$). These values were chosen to ensure atom loss of approximately 50%. After the spectroscopy pulse was completed, the atoms were released from the ODT and we performed absorption imaging on the $^1S_0 \rightarrow ^1P_1$ transition to measure atom loss as a function of the spectroscopy laser frequency.

For all four isotopes, data was taken across several days and referenced to the frequency comb. Then the frequencies were averaged to obtain a single line center for each isotope. A final isotope shift was found by subtracting the measured absolute frequencies relative to ^{88}Sr , and the total errors were added in quadrature. For the ^{87}Sr isotope shift, we weight the measurements of each excited-state hyperfine manifold $F' \in \{11/2, 9/2, 7/2\}$ to find the nominally unshifted line center in the absence of the hyperfine interaction [94]. However, it is important to note that this model fits three parameters (hyperfine A and B coefficients and an unshifted line center) from three isotope shifts, and thus is completely determined by the available data¹. A more accurate theory of higher order shifts from other fine structure levels will be necessary to assign a more accurate isotope shift for ^{87}Sr . This is currently an area of ongoing theoretical research [96].

To calculate the final value for the isotope shift, we also evaluated system-

¹We determined $|A| = 260085 \pm 2 \text{ kHz}$ and $|B| = 35667 \pm 21 \text{ kHz}$, consistent with previous results [95].

atic effects, as summarized in Table 5.1. Since many of the systematic effects are common to both isotopes, and the isotope shift is found from a difference in those frequencies, many potential systematic effects are common mode and cancel to a high degree. This is particularly true for the even isotopes, where there is no hyperfine structure. For example, even though a magnetic field is applied during the spectroscopy pulse for the even isotopes, the Zeeman shift is identical to within our experimental uncertainties, and does not lead to a correction to the final isotope shift. Therefore, as shown in Table 5.1, the remaining systematic effects are those that are not common mode: the density shift and recoil shift.

The density shift arises due to the different scattering lengths and atom numbers between different isotopes in our experiment. The cumulative effect is a non-zero differential density shift to the final isotope shift value. We experimentally determined this density shift for each isotope by measuring the line center at different atom numbers while keeping all other parameters the same. A linear fit allowed us to extrapolate from our operating atom number to a nominal “zero-density” frequency, yielding the systematic density shift shown in Table 5.1. The photon recoil shift [97] was also accounted for and was calculated from known physical quantities.

To first order, the $^1S_0 \rightarrow ^3P_1$ transition is magnetic field insensitive, and our measurements were performed at a low magnetic field of 0.05 mT (0.5 G) for the bosons and zero magnetic field for the fermion. Therefore both the first and second order Zeeman shifts were negligible at our level of accuracy. The stroboscopic procedure described above removed any AC Stark shifts due to the 1064-nm trapping beam. Finally, since the intensity in the 689-nm spectroscopy pulse was low (at most

0.1 mW/cm²) and the probe times were short (a few ms), systematic shifts from the probe pulse were below our experimental uncertainty.

Isotope Shift (kHz)	$^1S_0 \rightarrow ^3P_1$	$^1S_0 \rightarrow ^3P_0$
88-84	$351495.8 \pm 0.3 \pm 4.6$	$349656 \pm 1 \pm 10$
88-86	$163818.7 \pm 0.3 \pm 3.8$	$162939 \pm 2 \pm 11$
88-87	$62186.5 \pm 0.6 \pm 11.7$	$62171 \pm 1 \pm 23$
88-87 ($F' = 7/2$)	$-1351933.1 \pm 2.1 \pm 14.6$	
88-87 ($F' = 9/2$)	$-221676.6 \pm 0.4 \pm 26.9$	
88-87 ($F' = 11/2$)	$1241485.8 \pm 0.3 \pm 15.7$	

Table 5.2: Measured isotope shifts relative to ^{88}Sr . For ^{87}Sr ($^1S_0 \rightarrow ^3P_1$), contributions from the three excited-state hyperfine manifolds are weighted to establish the fine-structure line center. Uncertainties are one standard deviation and indicate statistical and systematic uncertainties.

After applying corrections for the systematic effects, the final values for $^1S_0 \rightarrow ^3P_1$ isotope shifts are shown in Table 5.2. The total systematic uncertainties are determined by adding the individual systematic uncertainties for each isotope in Table 5.1 in quadrature. Our results are consistent with a previous measurement of the ^{88}Sr - ^{86}Sr isotope shift, which reported a value of 163817.4 ± 0.2 kHz [98].

Systematic Shift (kHz)	88-87		88-86		88-84	
	88	87	88	86	88	84
Density	0.8 ± 1.6	-3.8 ± 1.2	0.2 ± 0.3	-0.9 ± 0.8	0.4 ± 0.9	-2.3 ± 0.9
Recoil	$4.7 \pm (< 0.1)$	$4.7 \pm (< 0.1)$	$4.7 \pm (< 0.1)$	$4.8 \pm (< 0.1)$	$4.7 \pm (< 0.1)$	$4.9 \pm (< 0.1)$
AC Stark	51 ± 5	42 ± 22	53 ± 5	51 ± 5	50 ± 5	51 ± 5
Thermal	-22 ± 4	-17 ± 4	-22 ± 4	-8 ± 4	-22 ± 4	-21 ± 4
2nd Order Zeeman	$-2.8 \pm (< 0.1)$	$0.0 \pm (< 0.1)$	$-2.8 \pm (< 0.1)$	$-2.8 \pm (< 0.1)$	$-2.8 \pm (< 0.1)$	$-9.1 \pm (< 0.1)$
Probe Power	3.5 ± 1.6	1.3 ± 0.3	3.5 ± 1.6	3.5 ± 1.6	3.6 ± 1.6	3.6 ± 1.6
Probe Duration	3.4 ± 3.3	3.3 ± 1.3	3.4 ± 3.3	3.4 ± 3.3	3.4 ± 3.3	3.4 ± 3.3
Total	39 ± 8	31 ± 22	40 ± 7	51 ± 7	37 ± 7	31 ± 7

Table 5.3: Systematic frequency shifts and one standard deviation uncertainties for the $^1S_0 \rightarrow ^3P_0$ transition. The three columns for ^{88}Sr correspond to three independent isotope shift measurements. Uncertainties indicate one standard deviation.

5.1.5 Measurement of the $^1S_0 \rightarrow ^3P_0$ isotope shifts

The procedure for measuring the 698-nm transition differed from the measurement of the $^1S_0 \rightarrow ^3P_1$ intercombination-line transition in several key ways. Since the clock transition is strictly forbidden by angular momentum considerations for the bosonic isotopes, a much larger field was necessary to induce a transition in these isotopes. For ^{88}Sr and ^{86}Sr a magnetic field of 10.96 ± 0.02 mT (109.6 ± 0.2 G) was used, and 19.79 ± 0.05 mT (197.9 ± 0.5 G) was used for ^{84}Sr . For measurements of ^{87}Sr , which is weakly allowed due to hyperfine mixing, we applied zero magnetic field. For all isotopes, the 698-nm spectroscopy pulse was applied for 2 s with typical peak intensities of 0.87 W/cm² for the even isotopes and 0.12 W/cm² for the odd isotope ($I_{\text{sat}} \approx 0.4$ pW/cm²). These values were chosen to ensure approximately 50% atom loss. Atom loss was induced by the light scattering and subsequent recoil of the 698-nm light, which ejected the atoms out of the trap. Representative line shapes for the $^1S_0 \rightarrow ^3P_0$ transitions are shown in Fig. 5.1 for each isotope.

The spectroscopy beam was aligned in the horizontal plane at an angle of approximately 45° with respect to the ODT, and was focused onto the atoms with a $1/e^2$ waist of 330 μm in the horizontal direction and 460 μm in the vertical direction. The beam was linearly polarized parallel to the magnetic field. Finally, because of the long interrogation time needed for sufficient atom loss (and therefore sufficient signal to noise), we were unable to apply the stroboscopic procedure used to measure the $^1S_0 \rightarrow ^3P_1$ transitions, resulting in large AC Stark shifts from the trapping beam. Due to the modified experimental procedure for the clock transition, addi-

tional systematic shifts included: thermal shifts, second order Zeeman shifts, and spectroscopy pulse shifts.

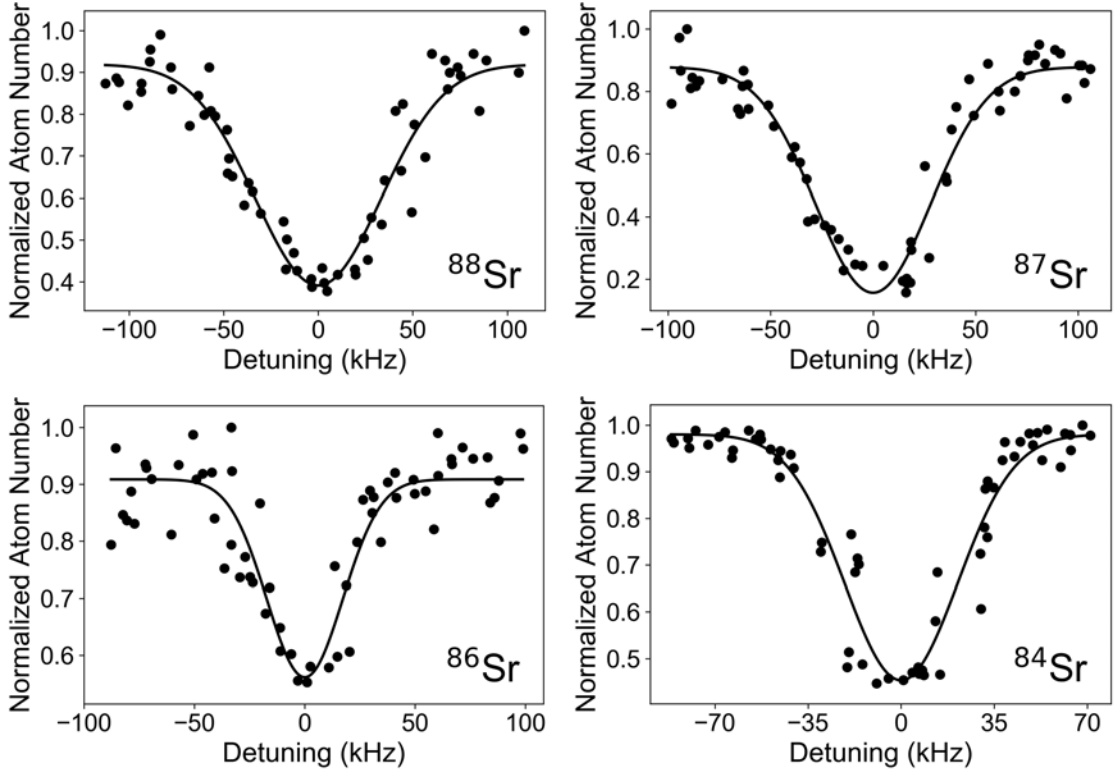


Figure 5.1: Spectroscopy of the $^1S_0 \rightarrow ^3P_0$ transition for each strontium isotope. The normalized atom number is shown as a function of the laser detuning. The solid line is a Gaussian fit to the data.

For the $^1S_0 \rightarrow ^3P_0$ transition, the dominant systematic effects were the AC Stark shift and what we call the thermal shift. The AC Stark shift arises from the differential polarizability of the 1S_0 and 3P_0 states at 1064 nm. The thermal shift arises from the inhomogeneous broadening and shift from the thermal motion in the intensity distribution of the ODT. Experimentally, the AC Stark shift was determined by measuring the resonance frequency as a function of the ODT intensity. However, varying the intensity of the ODT also varied the trap depth, which in turn varied the temperature of the atomic cloud. This led to additional shifts

in the resonance frequency due to both the Doppler shift and the inhomogeneous differential AC Stark shift. To distinguish the effects of the AC Stark shifts from the thermal shifts we took the thermal average using the Maxwell-Boltzmann distribution and modeled the scattering process from the spectroscopy pulse [92, 99] (see the Appendix 5.1.9). As shown in Table 5.3, the experimentally determined values for the AC Stark shift for each isotope agree with each other within the uncertainty. Therefore the AC stark shift is common mode and cancels to a high degree.

For even isotopes, the systematic shift for the first order Zeeman effect is zero since we probe a $J = 0 \rightarrow J' = 0$ transition with no hyperfine structure. To determine the second order Zeeman shifts for the $^1S_0 \rightarrow ^3P_0$ transitions we used our calibrated magnetic field measurements and the known second order Zeeman shifts [9], which are identical for all even isotopes. Spectroscopy of ^{87}Sr was performed at zero magnetic field, and so the Zeeman shift was well below other systematic effects [2].

The last systematic effects evaluated for the clock transition were related to the spectroscopy laser, occurring due to the relatively long probe time (2 s) and high peak intensities (0.87 W/cm²). To measure these systematics, the transition frequency was measured as a function of both pulse power and duration, and the shift was extrapolated to zero. Finally, the density shift and recoil shift were obtained using the same procedure as described for the $^1S_0 \rightarrow ^3P_1$ transition.

The final values for the isotope shift of the clock transition, including systematic corrections, are shown in Table 5.2. The systematic shifts are summarized in Table 5.3. The total systematic uncertainties for the clock transitions in Table 5.2 are determined by adding the individual systematic uncertainties for each isotope in

Table 5.3 in quadrature. Comparing to prior measurements of the ^{88}Sr - ^{87}Sr isotope shift which were all approximately $62188 \pm (< 1)$ kHz [82, 83, 85], our result of 62171 ± 24 kHz is consistent to well within one standard deviation.

5.1.6 King plot analysis

We performed a King plot analysis using our measured values of the isotope shifts, including the first measurements of the ^{88}Sr - ^{86}Sr and ^{88}Sr - ^{84}Sr isotope shifts for the clock transition. A King plot analysis is a systematic approach to quantitatively and visually analyze isotope shifts of different atomic transitions referenced to the same isotope by relating the isotope shifts between different transitions [25]. This is a function of the mass and field shift constants, which are independent of the isotopes and depend only on the transitions under consideration [100]. Specifically, the isotope shifts between isotopes of mass numbers A and A' on two transitions i and j can be written

$$\mu_{A,A'}\delta\nu_i^{A,A'} = K_i - \frac{F_i}{F_j}K_j + \frac{F_i}{F_j}\mu_{A,A'}\delta\nu_j^{A,A'}, \quad (5.1)$$

where $1/\mu_{A,A'} = 1/m_{A'} - 1/m_A$ is the inverse mass constant, m_A is the mass of isotope A [101], K_i is a constant associated with the mass shift of transition i , F_i is the field shift constant for transition i , and $\delta\nu_i^{A,A'} = \nu_i^{A'} - \nu_i^A$ is the isotope shift between isotopes A and A' on transition i [25, 75]. For our particular analysis, we have $A = 88$, and $A' \in \{87, 86, 84\}$, $i \equiv {}^1S_0 \rightarrow {}^3P_0$ at 698 nm, and $j \equiv {}^1S_0 \rightarrow {}^3P_1$ at 689 nm. An important point to note is that Eq. 5.1 describes a linear relationship

between isotope shifts of different transitions.

The King plot for our measured isotope shifts is shown in Fig. 5.2. A linear fit to all three points weighted by their uncertainties leads to a field shift constant ratio of $F_{698}/F_{689} = 0.987 \pm 0.008$ and $K_{698} - \frac{F_{698}}{F_{689}}K_{689} = 5.20 \pm 5.31$ GHz·amu, where the statistical and systematic uncertainties are added in quadrature. We have also performed a linear fit by replacing our measurement of the ^{88}Sr - ^{87}Sr 698-nm transition isotope shift with the more precise value from Ref. [85]. This leads to values of $F_{698}/F_{689} = 0.981 \pm 0.005$ and $K_{698} - \frac{F_{698}}{F_{689}}K_{689} = 8.56 \pm 3.45$ GHz·amu, which are consistent with values obtained using our measurement of the ^{88}Sr - ^{87}Sr 698-nm transition isotope shift.

Since there is some uncertainty in deriving the frequency for ^{87}Sr due to the hyperfine structure, we also fit the data after excluding this point to obtain a field shift constant ratio of $F_{698}/F_{689} = 0.998 \pm 0.002$ and $K_{698} - \frac{F_{698}}{F_{689}}K_{689} = -1.87 \pm 1.03$ GHz·amu where the uncertainties are propagated from the uncertainties of each point for both axes. Compared to this two-point linear fit, the ^{88}Sr - ^{87}Sr 689-nm isotope shift we determined would have to increase by 136.2 kHz to become consistent with a linear King plot. Given that our data points with their uncertainties lie well outside of the straight line fit to all three points, the results in Fig. 5.2 suggest a possible nonlinear contribution to Eq. 5.1, or may indicate significant uncertainties in the determination of the center-of-mass of the ^{87}Sr 3P_1 hyperfine structure. In particular, our data indicates a nonlinearity using the nonlinearity measure defined in Ref. [26]. Future theoretical and experimental studies should help to explain our observations, including better calculations of the hyperfine mixing within the 3P

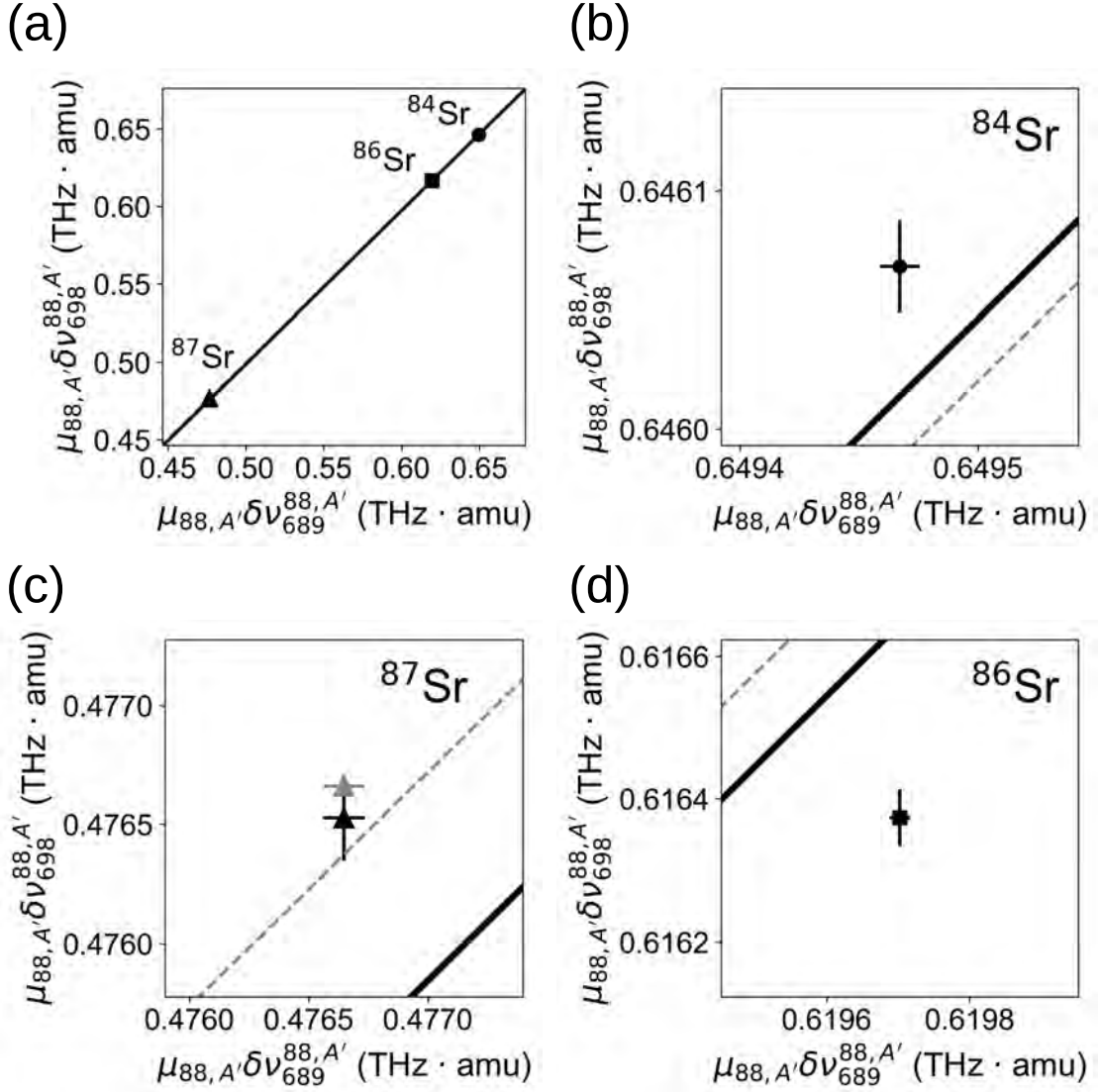


Figure 5.2: King plot of the measured strontium isotope shifts. (a) Linear fit to the three points derived from the six isotope shift measurements. Solid black line is a fit using all six of our measured isotope shifts, and the dashed gray line is a fit by replacing our measured ^{88}Sr - ^{87}Sr 698-nm transition isotope shift with the value from Ref. [85], which is more precise than our measurement. The black points are derived from our measurements, and the gray point is using the ^{88}Sr - ^{87}Sr 698-nm transition isotope shift from Ref. [85]. The fits are weighted by the uncertainties of each point. Error bars and the difference between the ^{87}Sr points derived from our measurement and from Ref. [85] are not visible at this scale. (b)-(d) Close up of each point in (a) with error bars shown.

states and a prediction of the King plot slope.

5.1.7 Conclusions

In summary, we have presented the first spectroscopy of the $^1S_0 \rightarrow ^3P_0$ clock transition in ^{86}Sr and ^{84}Sr , and reported their isotope shifts relative to ^{88}Sr . In conjunction with improved measurements of the intercombination line isotope shifts, we performed a King plot analysis and extracted constants related to the field and mass shifts. Hyperfine effects in ^{87}Sr complicate this analysis, but the experimental precision permitted by these two narrow optical transitions make it a rich testbed to benchmark state-of-the-art theory. Furthermore, it has been suggested that a comparison of isotope shifts between neutral and ionic strontium could set stringent limits on new physics [26, 77]. However, an improved theory, accounting for our observed nonlinearity would be essential. Alternatively, one could also perform this measurement with the radioactive bosonic isotope ^{90}Sr (half-life of approximately 29 years [102]) to avoid complications due to the hyperfine structure.

Future improvements on the measured frequencies will be possible by applying techniques successfully used with state-of-the-art strontium optical clocks, such as the use of magic-wavelength dipole traps to minimize the differential AC stark shift [103, 104] and optical lattices to suppress motional broadening and recoil shifts [97]. These advances should further suppress statistical and systematic errors on both transitions, allowing measurements with fractional uncertainties down to the level of 10^{-18} [11, 12, 81]. Our results, combined with other recent measurements

of isotope shifts in Ca^+ [105] and Sr^+ [106], will further help to refine atomic structure calculations and constrain new physics.

5.1.8 Acknowledgments

We thank Luis Orozco, Marianna Safronova, and Charles Clark for fruitful discussions and Nicholas Mennona for experimental assistance. This work was partially supported by the U.S. Office of Naval Research and the NSF through the Physics Frontier Center at the Joint Quantum Institute.

5.1.9 Appendix: Modeling the inhomogeneous broadening of the clock transition

In general, the AC Stark shift is different for different atomic states due to state-dependent polarizabilities. The exception to this is if one operates the dipole trap at specific laser wavelengths typically referred to as the “magic wavelength” where the ground and excited states experience the same AC Stark shifts. For strontium atoms, the magic wavelength is 813 nm for the 698-nm clock transition and 914 nm for the 689-nm intercombination transition [8, 11]. In our experiment, the optical dipole trap uses 1064-nm laser light, a wavelength where the two states, 1S_0 and 3P_0 , have different polarizabilities. This leads to inhomogeneous broadening which must be accounted for. The resulting lineshape is further complicated by the temperature of our atomic samples. Here we describe our method for modeling and accounting for this inhomogeneous broadening due to both the differential AC Stark

shift and the thermal shift.

We model the inhomogeneous broadening process using a semi-classical treatment of atom loss from the trap due to the spectroscopy pulse [99]. We can model the atom loss from the spectroscopy pulse after some probe time, by calculating the loss rate coefficient K . The time-dependent atom number in the presence of the spectroscopy pulse is governed by the differential equation

$$\frac{dN}{dt} = -K(\delta\omega, I, T, U_{\text{trap}})N, \quad (5.2)$$

where the loss rate coefficient K is a function of the laser detuning $\delta\omega = \omega_{\text{laser}} - \omega_0$ (ω_{laser} is the frequency of the probe laser and ω_0 is the bare atomic resonance frequency), the probe laser intensity I , the atomic cloud temperature T , and the dipole trap potential U_{trap} . The loss rate is modeled to be proportional to an ensemble average of the scattering rate over all atoms in the trap.

The scattering rate can be written [37]

$$\Gamma_{\text{scat}} = \frac{\Gamma}{2} \left(\frac{s_0}{1 + s_0 + (2\Delta/\Gamma)^2} \right), \quad (5.3)$$

where Γ is the transition linewidth, Δ is the effective detuning from resonance, $s_0 \equiv I/I_{\text{sat}}$ is the on-resonance saturation parameter, I is the excitation laser intensity, and I_{sat} is the saturation intensity. We rearrange this expression, pulling out

constant terms to write

$$\Gamma_{\text{scat}} \propto \frac{1}{(\Gamma'/2)^2 + \Delta^2}, \quad (5.4)$$

where $\Gamma' = \Gamma\sqrt{1 + s_0}$ is the saturation-broadened linewidth. For a thermal atom in a far-detuned optical dipole trap with a given phase space coordinate $(\mathbf{r}, \mathbf{p})$, Δ can be written

$$\Delta = \delta\omega - \frac{\mathbf{p} \cdot \mathbf{k}}{m} - (U_e(\mathbf{r}) - U_g(\mathbf{r})), \quad (5.5)$$

where the term $\delta\omega - \mathbf{p} \cdot \mathbf{k}/m$ is the Doppler-shifted laser frequency, $\mathbf{p}$ is the atomic momentum vector, $\mathbf{k}$ is the probe laser wavevector, m is the atomic mass, and $U_e(\mathbf{r}) - U_g(\mathbf{r})$ is the differential AC Stark shift which arises from different polarizabilities between the states e and g . Note that in the treatment here we neglect all other systematic frequency offsets which do not depend on position, since these appear simply as frequency offsets and do not cause any inhomogeneous effects. We also neglect gravity in our model since the atoms are tightly confined in this direction. We can approximate the trapping potential for the far-detuned optical trap as a

parabola and write

$$U_e(\mathbf{r}) - U_g(\mathbf{r}) = U_{e,0} + \frac{1}{2}m\bar{\omega}_e^2\mathbf{r}^2 - U_{g,0} - \frac{1}{2}m\bar{\omega}_g^2\mathbf{r}^2 \quad (5.6)$$

$$= U_{e,0} - U_{g,0} + \frac{1}{2}m(\bar{\omega}_e^2 - \bar{\omega}_g^2)\mathbf{r}^2 \quad (5.7)$$

$$= U_{e,0} - U_{g,0} + \frac{1}{2}m\bar{\omega}_g^2\mathbf{r}^2 \left(\frac{\bar{\omega}_e^2}{\bar{\omega}_g^2} - 1 \right) \quad (5.8)$$

$$= \Delta U_0 + \Delta U_{\text{trap}}(\mathbf{r}), \quad (5.9)$$

where $\bar{\omega}_g(\bar{\omega}_e)$ is the geometric mean of the ground (excited) state trap frequencies in all three dimensions and the trap is effectively spherical in these coordinates. Since $\bar{\omega}_i \propto \sqrt{\alpha_i}$, where α_i is the AC polarizability of state $i \in \{g, e\}$, we find

$$\Delta U_{\text{trap}}(\mathbf{r}) = U_{\text{trap}}(\mathbf{r}) \left(\frac{\alpha_e}{\alpha_g} - 1 \right), \quad (5.10)$$

where $U_{\text{trap}} = m\bar{\omega}_g^2\mathbf{r}^2/2$. For 1064-nm light, with g the 1S_0 state and e the 3P_0 state, we compute $\alpha_e/\alpha_g \approx 0.7$. This can also be written as a re-scaling of the trap potential, such that

$$\Delta U_{\text{trap}}(\mathbf{r}) = \alpha U_{\text{trap}}(\mathbf{r}), \quad (5.11)$$

with $\alpha = (\alpha_e/\alpha_g - 1) \approx -0.295$. Note that operating the dipole trap at the magic wavelength would lead to $\alpha_e = \alpha_g$, which means $\alpha = 0$, and therefore the spatial dependence would drop out of Eq. 5.5.

We now turn our attention to solving for the loss rate coefficient K by taking

an ensemble average over the scattering rate expressed in Eq. 5.4 using the detuning defined in Eq. 5.5. Because we are interested in deriving a lineshape function which can be fit to experimentally measured atom loss data, we ignore normalization and overall constant terms which can be condensed into a single fit parameter. Taking the ensemble average of Eq. 5.4 leads to

$$K \propto \int d^3\mathbf{r} e^{-U_{\text{trap}}(\mathbf{r})/(k_B T)} \int d^3\mathbf{p} e^{-\mathbf{p}^2/(2mk_B T)} \left[\frac{1}{(\Gamma'/2)^2 + (\delta\omega - \Delta U_0 - \mathbf{p} \cdot \mathbf{k}/m - \alpha U_{\text{trap}}(\mathbf{r}))^2} \right], \quad (5.12)$$

where we have taken an integral over phase space $(\mathbf{r}, \mathbf{p})$ weighted by the Boltzmann factor. Here, k_B is the Boltzmann constant.

We wish to make this dimensionless to easily work in a numerical fitting routine with experimental data. Focusing on the integral $d^3\mathbf{p} = dp_x dp_y dp_z$ first, we can choose $\hat{p}_z$ to point along $\mathbf{k}$. Thus, $\mathbf{p} \cdot \mathbf{k} = p_z k$, and the Boltzmann factor can be rewritten

$$e^{-\mathbf{p}^2/(2mk_B T)} = e^{-(p_x^2 + p_y^2)/(2mk_B T)} e^{-p_z^2/(2mk_B T)} \quad (5.13)$$

The integral over p_x and p_y now factors out, and can be brought into an overall scale factor. We define the dimensionless variable $y \equiv p_z/\sqrt{2mk_B T}$. After defining $\beta \equiv k\sqrt{2k_B T/m}$, this becomes $p_z k/m = \beta y$. In convenient units, for ^{88}Sr and $2\pi/k = 698 \text{ nm}$, we get $\beta/2\pi = 19.7 \text{ kHz} \cdot \sqrt{T}$ with T measured in μK . This parameterization of y serves to scale the momentum p_z to the most probable momentum at a given

temperature.

Putting it all together, the integral from Eq. 5.12 becomes

$$K \propto \int d^3\mathbf{r} e^{-U_{\text{trap}}(\mathbf{r})/(k_B T)} \int dy e^{-y^2} \left[\frac{1}{(\Gamma'/2)^2 + (\delta\omega - \Delta U_0 - \beta y - \alpha U_{\text{trap}}(\mathbf{r}))^2} \right]. \quad (5.14)$$

With regards to the integral over $\mathbf{r}$, since we have scaled the trap to be effectively spherical, we can write $U_{\text{trap}}(\mathbf{r}) = f(r^2)$. Thus, we can pull the angular integral from $d^3\mathbf{r} \equiv r^2 \sin\theta dr d\theta d\phi$ into an overall constant, leaving just the integral in r given by

$$K \propto \int r^2 dr e^{-r^2 m \bar{\omega}_g^2 / (2k_B T)} \int dy e^{-y^2} \left[\frac{1}{(\Gamma'/2)^2 + (\delta\omega - \Delta U_0 - \beta y - \alpha m \bar{\omega}_g^2 r^2 / 2)^2} \right], \quad (5.15)$$

where we replaced $U_{\text{trap}}(\mathbf{r})$ with its explicit form $m \bar{\omega}_g^2 r^2 / 2$.

Defining the dimensionless variable $x \equiv r \sqrt{m \bar{\omega}_g^2 / (2k_B T)} = r (\bar{\omega}_g k / \beta)$, which scales r by the ratio of the trap potential energy to the thermal energy $k_B T$, we can rewrite the integral as

$$K \propto \int dx x^2 e^{-x^2} \int dy e^{-y^2} \left[\frac{1}{(\Gamma'/2)^2 + (\delta\omega - \Delta U_0 - \beta y - (\frac{\alpha m}{2k^2}) \beta^2 x^2)^2} \right], \quad (5.16)$$

where for our system, $\alpha m / (2k^2) \approx -2.52 \times 10^{-6} \text{ s}$.

Returning to Eq. 5.2, we use our expression Eq. 5.16 for K to solve for atom

number and obtain

$$\frac{N(\tau)}{N(0)} = e^{-K\tau}, \quad (5.17)$$

which can be used as an integral function to fit the four parameters $\{a, (\omega_0 + \Delta U_0(I_{\text{trap}})), \Gamma', \beta\}$, where a is an overall normalization factor for K , in a least-squares minimization routine. We keep the $\Delta U_0(I_{\text{trap}})$ term explicit and highlight its dependence on the optical dipole trap laser intensity I_{trap} . We use this expression to extract the AC Stark shift systematic correction.

Note that in theory, the integral in Eq. 5.16 ranges over the entire real line. In our numerical implementation, we truncate these integrals at finite values. In our experiment, we typically have $U_{\text{trap}} \sim 160$ kHz, and so $U_{\text{trap}}/(k_B T) \sim 8$ and we take the position integral out to five times the thermal energy scale. Since the integrand is convolved by a gaussian, continuing the integration further in the wings contributes only marginally to the final value, and the truncation does not change the result above other uncertainties. The ratio $U_{\text{trap}}/(k_B T) \sim 8$ also allows us to approximate the trap as harmonic.

As an example, we perform a fit using Eq. 5.17 to the loss spectra shown in Fig. 5.3. It is difficult to visually differentiate the quality of the fit between the full integral lineshape and a simple Gaussian model, but there is a non-negligible thermal line shift from a full accounting of the lineshape as is evident in the fit parameters. To account for this systematic shift, we numerically simulate the systematic Gaussian fit offset as a function of temperature, and find it to be -7.6 ± 0.3 kHz/ μ K, as shown

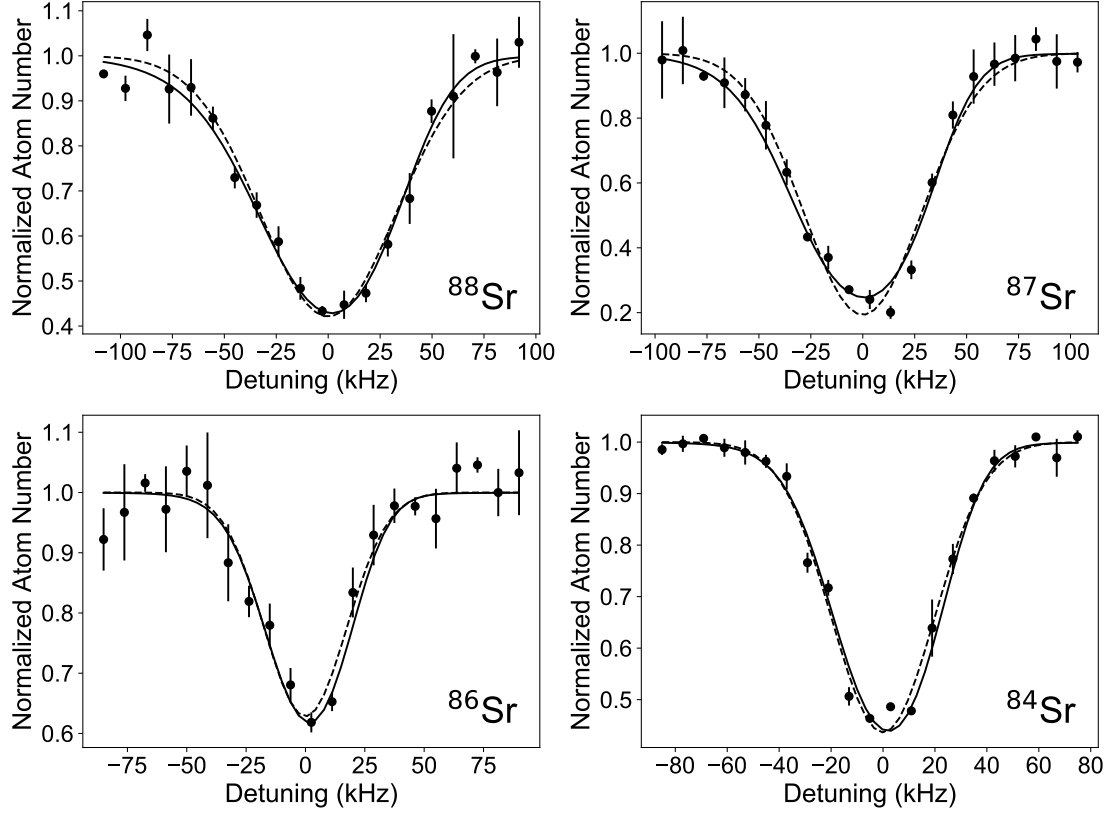


Figure 5.3: Lineshape curves for the 698-nm clock transition. The curves include a Gaussian model (dashed line) and a full lineshape model (solid line) fit to the averaged data points (circle points). In both cases, the fit error on the centroid is roughly 1 kHz, however the full lineshape model fits a different ω_0 which varies as a function of temperature and is red of the Gaussian line center by up to 20 kHz. This is attributable to the thermal distribution of atoms in a dipole trap with inhomogeneous AC Stark shifts.

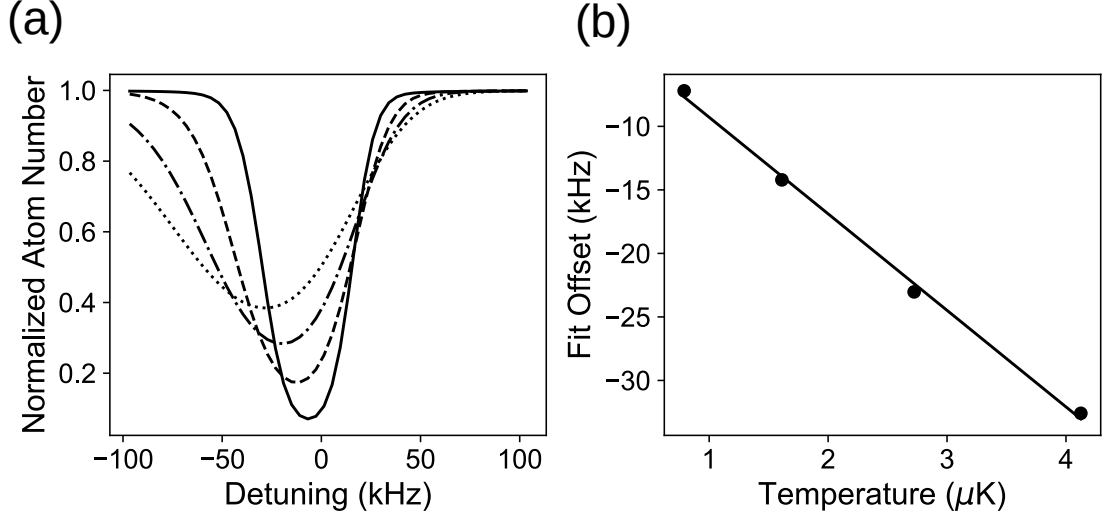


Figure 5.4: Effects of thermal line shift on the clock transition. (a) Lineshape simulations as a function of temperature with 0.79 μK (solid line), 1.6 μK (dashed line), 2.7 μK (dash-dotted line), and 4.1 μK (dotted line). (b) Systematic offset to the Gaussian fitted center as a function of temperature and a linear fit to the data extracted from the simulation in (a).

in Fig. 5.4. With this result and measured temperatures of {2.9 μK , 2.2 μK , 1.1 μK , 2.7 μK }, we obtain a systematic frequency shifts of $\{-22 \pm 4 \text{ kHz}, -17 \pm 4 \text{ kHz}, -8 \pm 4 \text{ kHz}, -21 \pm 4 \text{ kHz}\}$ for $\{^{88}\text{Sr}, ^{87}\text{Sr}, ^{86}\text{Sr}, ^{84}\text{Sr}\}$ respectively.

5.2 Publication: Confinement of an alkaline-earth element in a grating magneto-optical trap

Our work on the strontium grating MOT was completed in collaboration with the cold atom vacuum standards (CAVS) group at the National Institute of Standards and Technology (NIST), especially Daniel Barker, who suggested we try creating a strontium grating MOT. Our laboratory had the infrastructure at that point to run strontium laser cooling experiments, and the group at NIST had grating

chips, custom vacuum pieces, and experience with grating MOTs. Ananya Sitaram and I were the senior graduate students on the project together, and we were able to get our blue MOT up and running relatively quickly. The following publication, written up during the COVID-19 pandemic and published as Ref. [58], is the result of that work. We were able to create and characterize the first grating MOT for alkaline earth atoms.

5.2.1 Abstract

We demonstrate a compact magneto-optical trap (MOT) of alkaline-earth atoms using a nanofabricated diffraction grating chip. A single input laser beam, resonant with the broad $^1S_0 \rightarrow ^1P_1$ transition of strontium, forms the MOT in combination with three diffracted beams from the grating chip and a magnetic field produced by permanent magnets. A differential pumping tube limits the effect of the heated, effusive source on the background pressure in the trapping region. The system has a total volume of around 2.4 L. With our setup, we have trapped up to 5×10^6 ^{88}Sr atoms, at a temperature of approximately 6 mK, and with a trap lifetime of approximately 1 s. Our results will aid the effort to miniaturize quantum technologies based on alkaline-earth atoms.

5.2.2 Introduction

Laser-cooled alkaline-earth atoms have applications in a wide range of quantum devices, including atomic clocks [84, 107, 108], gravimeters [31], and spaceborne

gravitational wave detectors [33, 34]. The transition from a laboratory to field-based applications will require a drastic reduction in the size and complexity of laser-cooling systems. For example, proposals to detect gravitational waves using alkaline-earth atoms require atom interferometers capable of being installed in satellites [34]. Compact versions of these laser-cooled systems are also necessary in order make the unprecedented accuracy of alkaline-earth atomic clocks widely accessible [36].

Laser-cooling experiments typically use a magneto-optical trap (MOT) to capture, cool, and confine the atoms. Conventional MOTs use three orthogonal pairs of well-balanced, counterpropagating laser beams to confine atoms at the center of a quadrupole magnetic field. As such, MOTs require large vacuum chambers with optical access along all axes, and have many degrees of freedom in alignment and polarization. Compound optics can generate all necessary beams from a single input beam, reducing the complexity of the optical setup. For example, pyramidal retro-reflectors maintain the beam geometry of conventional MOTs [109]. However, the MOT forms inside the retro-reflecting optic, limiting optical access [110, 111]. Tetrahedral reflectors form the MOT above the optic, maintaining optical access but breaking the geometry of a conventional trap by using only four beams [112]. Tetrahedral MOTs can also be planarized by using diffraction gratings [46, 113, 114]. Thus far, only experiments with alkali atoms have been successfully miniaturized using such grating MOTs [47, 113, 114]. Here, we demonstrate a compact, grating MOT system for alkaline-earth atoms.

Alkaline-earth atoms pose unique challenges to miniaturization. First, sources for alkaline-earth atoms must be heated to high temperatures (over 350 °C) to

create sufficient flux of atoms to load a MOT. Outgassing from the hot source can increase the background pressure, decreasing the trap lifetime, and equilibrium atom number. Second, alkaline-earth atoms require large magnetic field gradients (on the order of 5 mT/cm), often created with large water cooled coils [115]. Third, with the high Doppler temperature of the broad $^1S_0 \rightarrow ^1P_1$ transition, and lack of sub-Doppler cooling, strontium and other alkaline-earth systems usually operate a second, subsequent MOT on the narrow $^1S_0 \rightarrow ^3P_1$ transition to achieve lower temperatures. The ideal compact system must have the capability to operate at the two different cooling wavelengths.

Our system, designed around a diffraction grating chip, mitigates the above issues associated with miniaturizing a MOT for alkaline-earth atoms (see Fig. 5.5). First, a 3 cm long differential pumping tube separates the vacuum chamber into two regions: the source chamber and the science chamber. The source chamber contains a vacuum pump and a low-outgassing dispenser [116] that vaporizes strontium atoms. The atoms then travel through the differential pumping tube before entering the science chamber. Second, we create the magnetic field gradient for the MOT using permanent magnets, which are less complex than typical, water-cooled coils. The magnetic field gradient extends into the differential pumping tube, forming an effective Zeeman slower when combined with the input laser beam. Lastly, the first order diffraction efficiency of the grating we use is optimal at a wavelength of 600 nm, a middle ground between the two laser-cooling wavelengths (461 nm and 689 nm) for strontium. Our compact alkaline-earth grating MOT system also maintains the optical access and achieves the atom number necessary for future quantum devices.

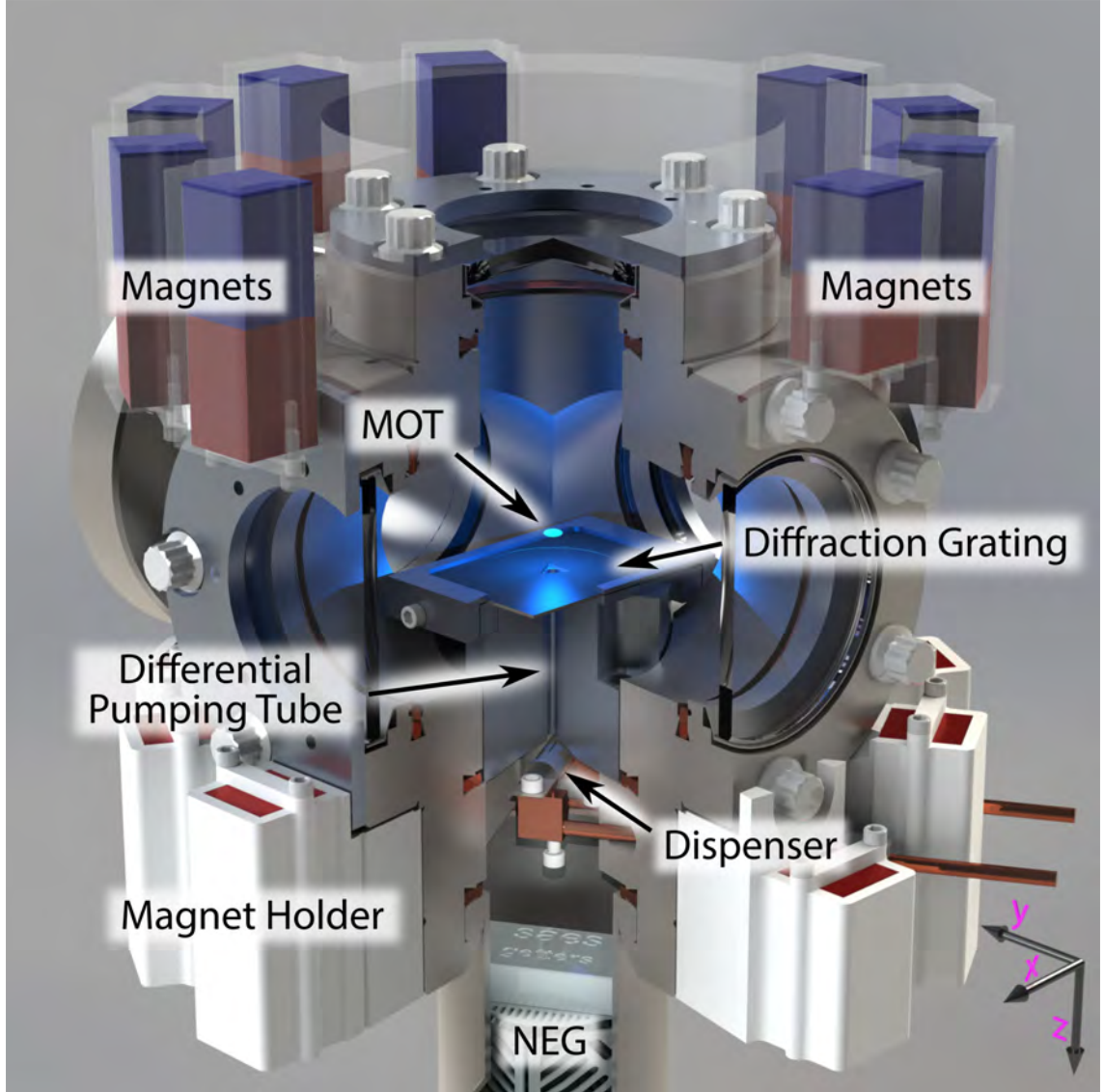


Figure 5.5: A cut-away model of the grating MOT system, with coordinates specified in the bottom right. The diffraction grating in the middle of the science chamber diffracts light to form the MOT beams. The input laser beam (not shown) propagates along the $+\hat{z}$ direction. 3D-printed magnet holders position permanent magnets around the top and bottom of the chamber. The top magnet holder is translucent to show the configuration of the magnets in the holders. The polarity of the magnets is shown with red and blue coloring. The dispenser source sits below the differential pumping tube and is pumped by a non-evaporable getter (NEG) pump.

5.2.3 Apparatus

Our apparatus is shown in Fig. 5.5. The vacuum system is comprised of two chambers, separated by a 3 cm long, 3 mm diameter differential pumping tube with an N_2 conductance of 0.11 L/s. The MOT is located in a science chamber with four CF275 [91] viewports, and pumped with a 75 L/s ion pump (not shown). The source chamber is located below the differential pumping tube and is pumped with a 40 L/s non-evaporable getter (NEG) pump. Our source of Sr atoms is a 3D-printed titanium dispenser, described in Ref. [116]. We run a current between 12 A and 14 A through the dispenser, effusing strontium towards the differential pumping tube. Together, the source and science chambers are approximately 2.4 L in volume, although this estimate does not include the ion pump or the magnet holders. In typical strontium experiments, the sources alone are often at least 2 L in volume. The base vacuum pressure of 2×10^{-7} Pa in the science chamber could be improved by replacing the large ion pump with a small hybrid NEG/ion pump, which would also reduce the size of the apparatus.

The grating chip is located above the differential pumping tube, and has a triangular hole through its center, allowing atoms to enter the science chamber. The grating chip was fabricated at the National Institute of Standards and Technology, and consists of three linear gratings arranged in a triangle. The parameters of the chip are the same as those in Ref. [47], except with a trench depth of 150(2) nm. This trench depth minimizes the 0th order diffraction at 600 nm, which is between the 461 nm and 689 nm cooling transition wavelengths for strontium. Each linear

grating diffracts 32(1) % of the normally-incident 461 nm light into each of the ± 1 diffraction orders with an angle of $27.0(5)^\circ$. 4 % of the light is diffracted into each of the ± 2 orders, and 11 % is diffracted into the 0 order, with the remainder lost due to the aluminum coating [117]. The 0 order light does not disrupt the MOT because of the hole in the grating directly beneath it. The diffraction angle is a trade-off between confinement of the 461 nm MOT and overlap with the second stage 689 nm MOT beams. For normally incident, circularly polarized light, the stokes parameters of the grating chip at 461 nm are $Q = -0.23(1)$, $U = -0.13(1)$, $V = 0.96(1)$, where $Q = 1$ ($Q = -1$) corresponds to s (p) polarization defined relative to the plane of reflection for each linear grating.

Two sets of grade N52 NdFeB magnets, arranged roughly in a dodecagon, create the magnetic field for the MOT. Within each set, the poles of the magnets are aligned. The magnets are housed in 3D-printed magnet holders made of polylactic acid (PLA) that are designed to produce a compact setup with high magnetic field gradients, as shown in Fig 5.5. Due to the geometric constraints of the vacuum chamber, the configuration of magnets is asymmetric, and the principal axes are rotated from those Fig 5.5. We achieve maximum gradients of $\{3.5 \text{ mT/cm}, 2.7 \text{ mT/cm}, 6.2 \text{ mT/cm}\}$ along the $\{\hat{x}', \hat{y}', \hat{z}\}$ axes, respectively, where $\hat{x}'$ and $\hat{y}'$ are rotated by $-\pi/6$ from $\hat{x}$ and $\hat{y}$. By removing magnets from the holders, we can lower the gradient to $\{1.9 \text{ mT/cm}, 1.9 \text{ mT/cm}, 3.8 \text{ mT/cm}\}$ along the $\{\hat{x}', \hat{y}', \hat{z}\}$ axes, respectively. The field gradient extends to $z \approx 50 \text{ mm}$, where $z = 0$ corresponds to the $\mathbf{B} = 0$ and $z \approx 40 \text{ mm}$ corresponds to the position of the source.

A single laser beam, red-detuned from the $^1S_0 \rightarrow ^1P_1$ transition at 461 nm,

enters through the top viewport along the $+\hat{z}$ axis and is normally incident upon the diffraction grating chip. The input MOT beam has a $1/e^2$ radius of 12 mm and a maximum power of 92 mW. For the $^1S_0 \rightarrow ^1P_1$ transition with natural linewidth $\Gamma/2\pi = 30.5$ MHz, $I_{\text{sat}} = 40.3$ mW/cm², giving a maximum peak $I/I_{\text{sat}} \approx 1$. Intensities I/I_{sat} reported herein always refer to the peak intensity of the input beam. The central portion of the beam continues through the hole in the diffraction grating and through the differential pumping tube. This beam, combined with the magnetic field gradient, allows for a small amount of initial slowing of the atoms, similar to a Zeeman slower. Atoms can be lost from the MOT because the excited 1P_1 state decays at a rate of 610 s^{-1} to the 1D_2 state, which in turn decays to the 3P manifold. To mitigate the atom loss, two repump lasers, with wavelengths 679 nm and 707 nm, address the $^3P_0 \rightarrow ^3S_1$ and $^3P_2 \rightarrow ^3S_1$ transitions, respectively. More information on the repump scheme can be found in Ref. [42]. The repump beams are combined together on a 50/50 beam splitter, and then combined with the input MOT beam using a polarizing beam splitter.

We use absorption and fluorescence imaging along $\hat{x}$ to characterize the MOT. Absorption images are taken after the MOT atom number equilibrates using a probe beam resonant with the $^1S_0 \rightarrow ^1P_1$ transition with $I/I_{\text{sat}} \approx 0.01$. We use the atom number from the absorption images to calibrate the atom number extracted from fluorescence images taken during loading. The Labscript suite software [67] controls the experiment and data collection. More detailed information on the laser systems can be found in Ref. [5].

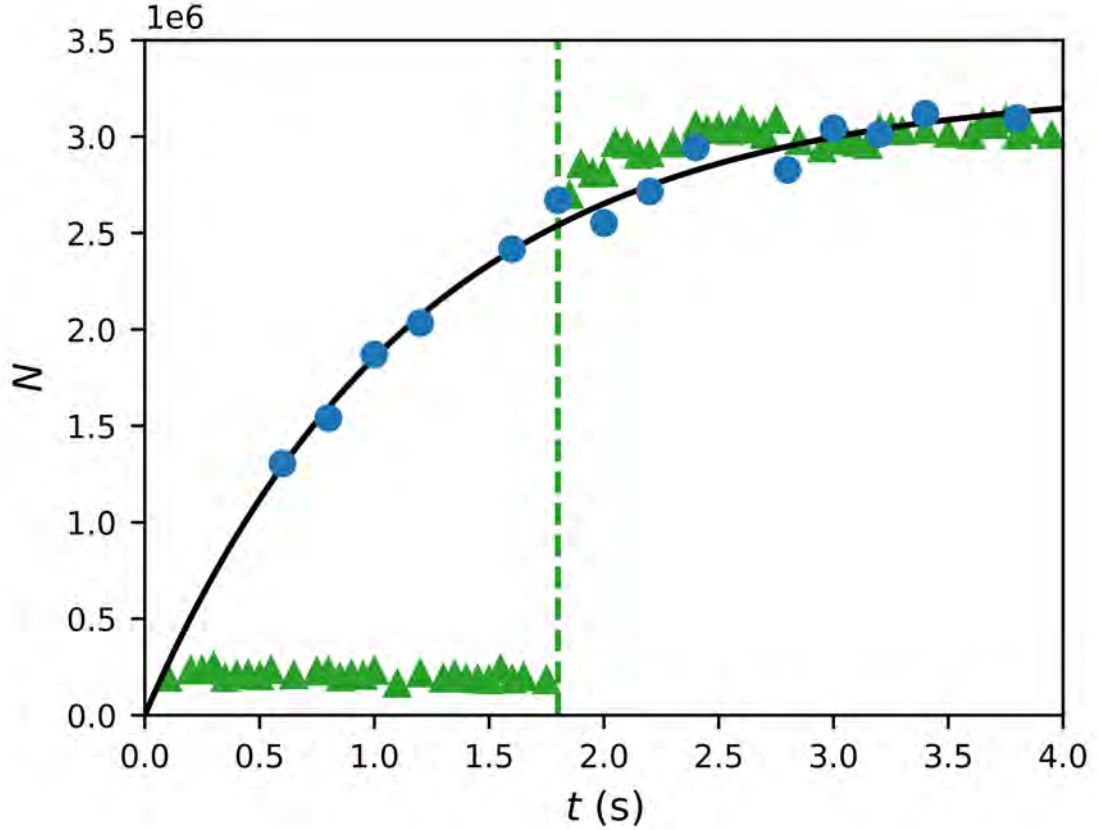


Figure 5.6: MOT loading curves with a source current of 13 A, $I/I_{\text{sat}} = 1$, axial magnetic field gradient of 6.2 mT/cm, and detuning $\Delta/\Gamma = -1$. The blue dots show the MOT atom number N as a function of time t . The black curve is a fit to Eq. 5.18. The green triangles show the MOT loading without repump lasers, and subsequent recapture from the metastable reservoir. The dashed line indicates when the repump lasers were turned on.

5.2.4 Results

We measure atom number, loading rate, lifetime, and temperature to characterize the MOT. During each experimental shot we take a sequence of fluorescence images while the MOT loads and construct a loading curve. Fig. 5.6 shows typical loading curves at an axial magnetic field gradient of 6.2 mT/cm. For a MOT with no light assisted collisions, the loading rate R , MOT lifetime τ , and equilibrium

atom number $N_0 = R\tau$, are extracted by fitting each loading curve to the single exponential

$$N(t) = R\tau(1 - e^{-t/\tau}). \quad (5.18)$$

An example fit is shown with a solid black curve in Fig. 5.6. The quality of the fit to Eq. (5.18) indicates light assisted collisions and secondary scattering are negligible. At the higher gradient of 6.2 mT/cm, we observe typical loading rates of $4 \times 10^6 \text{ s}^{-1}$ and a vacuum-limited lifetime of 1 s. We observe a similar loading curve at the lower gradient of 3.8 mT/cm.

Fig. 5.7 and Fig. 5.8 show the MOT parameters as a function of detuning from resonance and intensity, respectively. We find the maximum atom number of approximately 4×10^6 at a source current of 13 A and $\Delta/\Gamma \approx -1$, a typical detuning for a conventional 6-beam Sr MOT [118–120]. As shown in Fig. 5.8(a), the atom number continues to increase with I/I_{sat} , even at our maximum intensity, indicating that more laser power would be beneficial. The increase in $N_0 = R\tau$ is only partially due to the increase in the loading rate R , shown in Fig. 5.8(b). Part of the atom number increase is due to an increase in the lifetime with intensity, shown in Fig. 5.8(c). The lifetime increase suggests that the trap depth is increasing with laser power, which in turn increases the escape velocity for a Sr atom that undergoes a background gas collision [121, 122]. However, the interplay between MOT temperature and tighter radial confinement with increasing intensity may also play a role.

We can also use the MOT to continuously load a magnetic trap, which con-

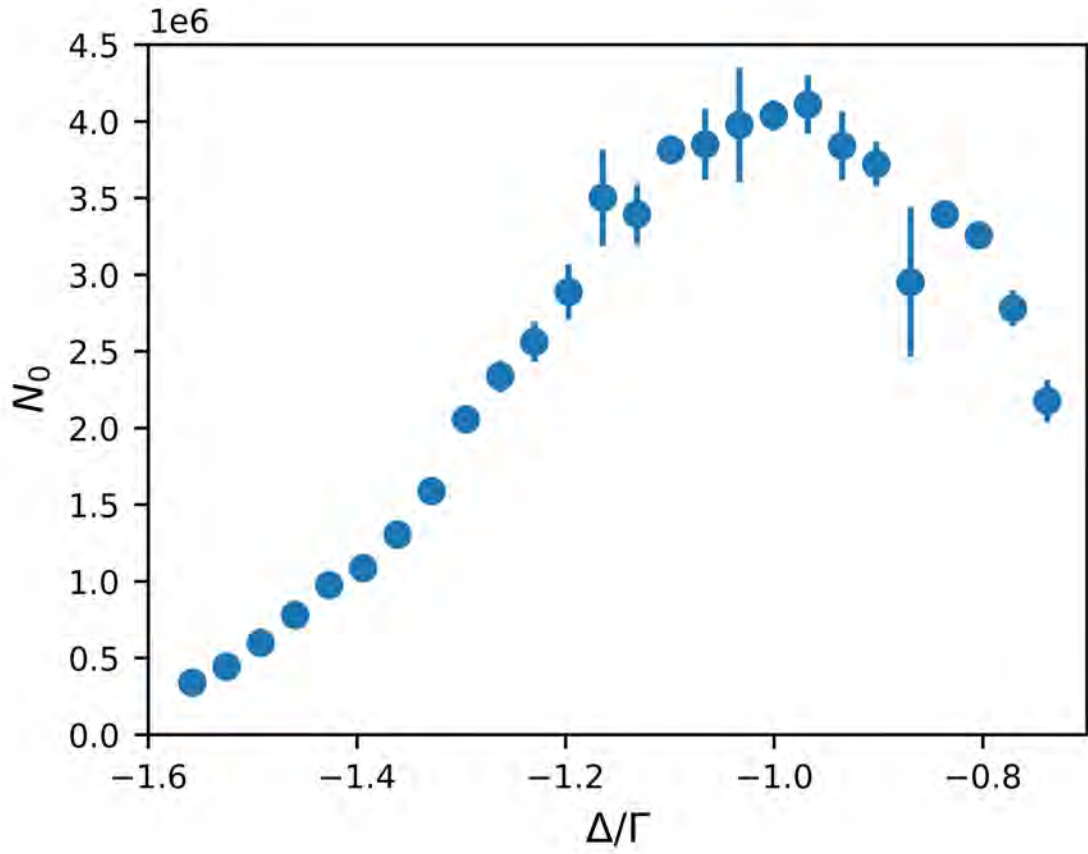


Figure 5.7: The equilibrium atom number N_0 as a function of MOT beam detuning Δ/Γ , with a source current of 13 A, axial magnetic field gradient of 6.2 mT/cm, and $I/I_{\text{sat}} = 1$. The optimal detuning is $\Delta/\Gamma = -1$. Most of the error bars are smaller than the data points, and represent the standard error about the mean.

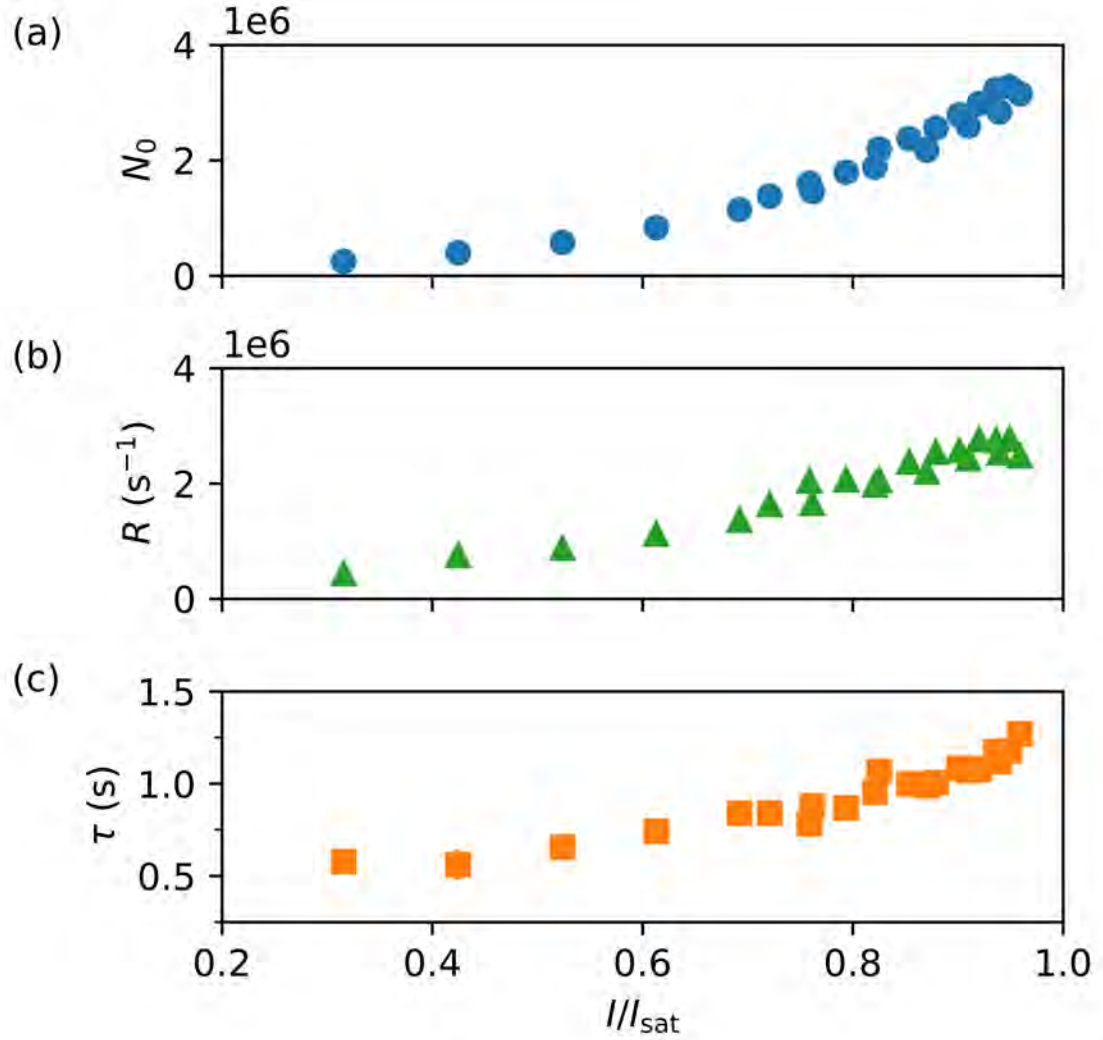


Figure 5.8: MOT loading parameters as a function of I/I_{sat} : (a) equilibrium atom number N_0 , (b) loading rate R , and (c) lifetime (τ). Here, the source current is 13 A, the detuning $\Delta/\Gamma = -1$, and axial field gradient is 6.2 mT/cm. The error bars on the points are comparable to the marker size, and represent the standard error about the mean.

sists of atoms that are trapped in the metastable $^3\text{P}_2$ state. With strontium, the metastable magnetic trap is often used to increase the capture of rare isotopes and was key to the realization of quantum degeneracy [119, 123]. By operating the MOT without repump light, atoms are shelved in the $^3\text{P}_2$ state where they are trapped by the MOT magnetic field. When the repump light is turned on after atoms have accumulated in the magnetic trap, we see a sharp increase in the MOT atom number as shown in Fig. 5.6. The recovery confirms that atoms are being caught and held in the magnetic trap, however we do not see a transient enhancement above the equilibrium atom number as demonstrated elsewhere [124]. Given our densities, and vacuum-limited atom number, we would not expect enhancement from magnetic trap loading. Adding a depumping laser could enhance the loading rate of the magnetic trap and increase the atom number [43].

We investigate the effect of the source current on the atom number, loading rate, and lifetime, shown in Fig. 5.9. The source current sets the temperature of the source, which in turn determines both the vapor pressure and the average velocity of atoms leaving the source. At our highest achievable source current of 14 A, limited by the ampacity of our electrical feedthroughs, we trap 5×10^6 atoms, but have still not saturated the atom number. Based on the fit presented in Ref. [116], we estimate the source temperature at 13 A to be over 600 °C. When the source current is increased from 0 A to 13 A, the vacuum pressure in the science chamber increases by 3×10^{-8} Pa, suggesting that the differential pumping is sufficient. The increase in pressure is consistent with the small lifetime decrease shown in Fig. 5.9(c).

To determine the temperature of the MOT, we measure the width of the atomic

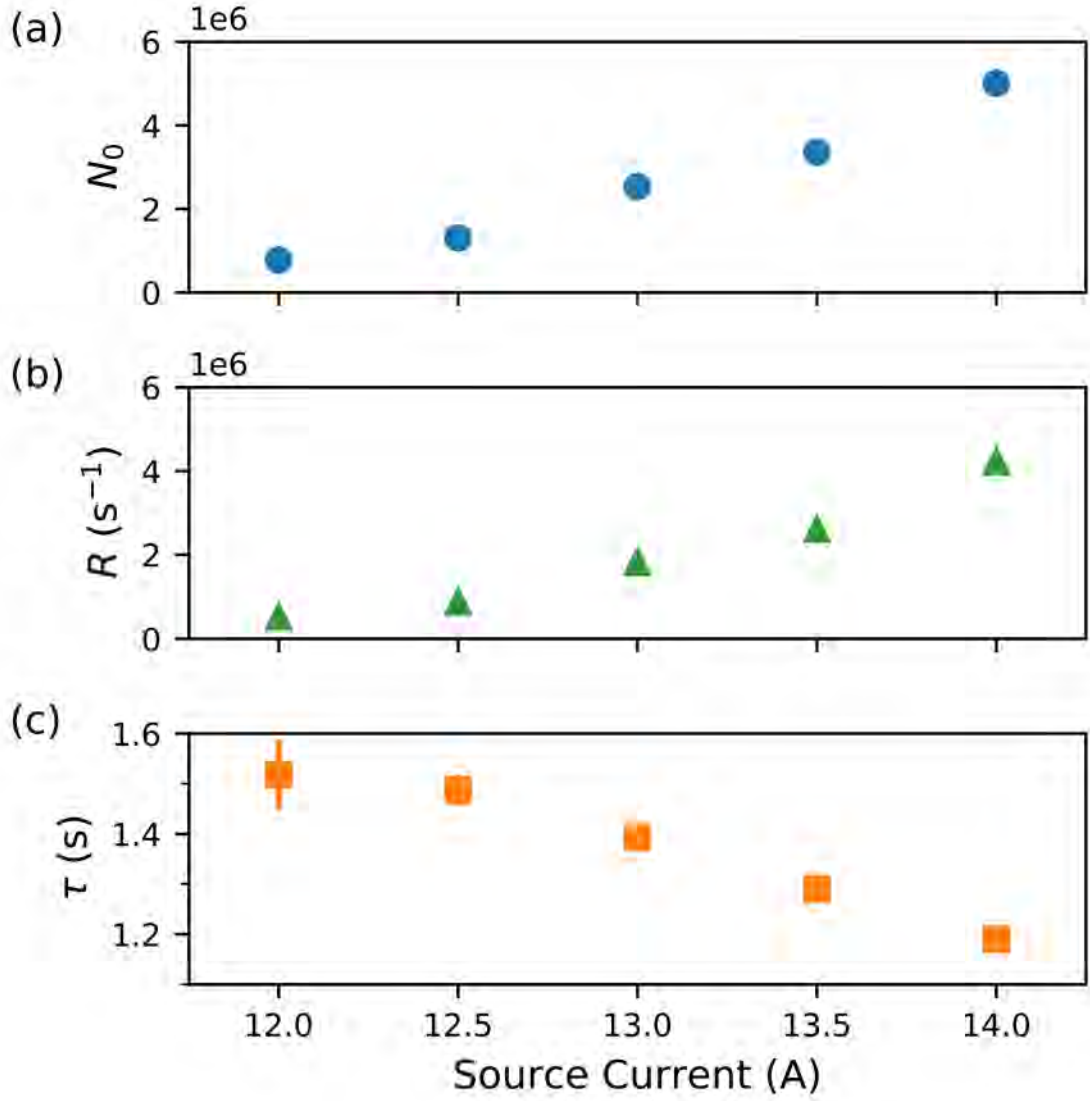


Figure 5.9: MOT loading parameters as a function of source current: (a) equilibrium atom number N_0 , (b) loading rate R , and (c) lifetime (τ). Here, $I/I_{\text{sat}} \approx 1$, $\Delta/\Gamma = -1$, and axial field gradient is 6.2 mT/cm. The error bars on most of the points are comparable to the marker size, and represent the standard error about the mean.

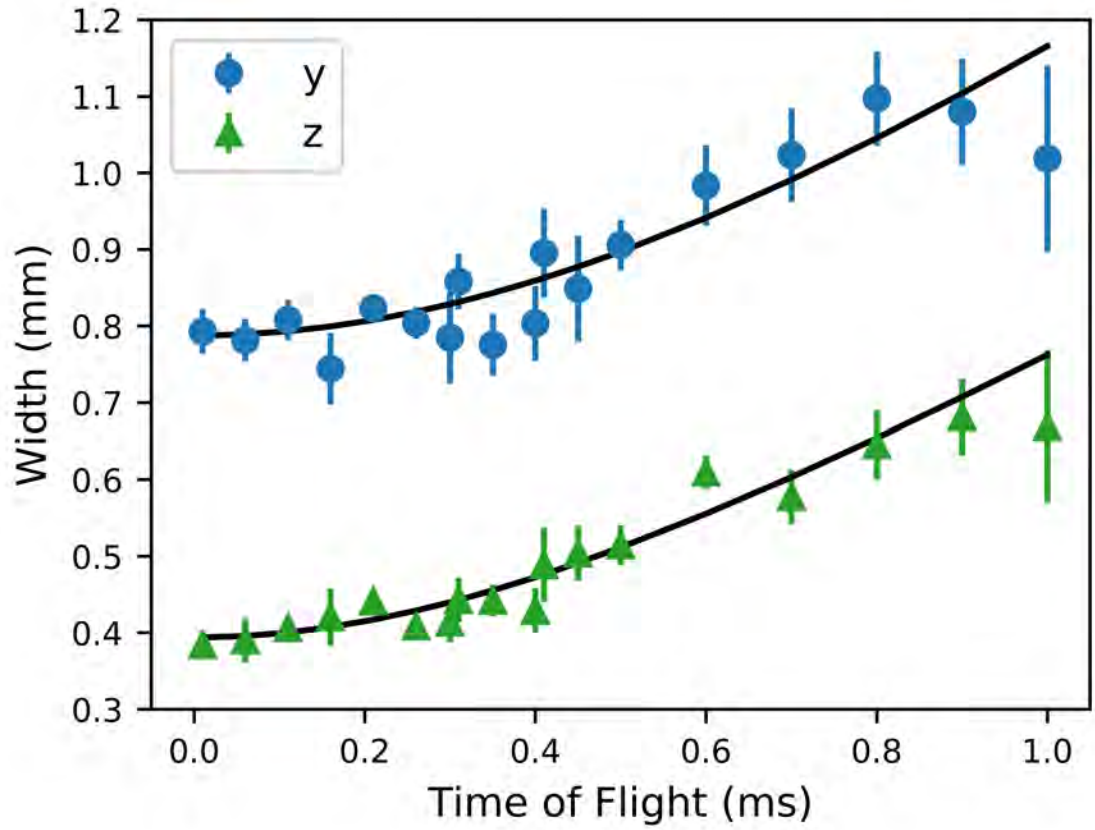


Figure 5.10: Temperature measurement of the atomic cloud. The rms width of the atomic cloud in the $\hat{y}$ direction (blue circles) and in the $\hat{z}$ direction (green triangles) are plotted against time of flight and fitted to the expansion function discussed in the text (black curves). The calculated temperatures based on the fits are 7.8(9) mK and 4.6(4) mK for $\hat{y}$ and $\hat{z}$, respectively, where the errors in parentheses are one standard deviation. This data was taken with a source current of 13 A, $I/I_{\text{sat}} = 1$, axial magnetic field gradient of 6.2 mT/cm, and detuning $\Delta/\Gamma = -1$. The error bars represent the standard error about the mean.

cloud as it expands in time of flight, shown in Fig. 5.10. A Gaussian fit extracts the root-mean-square (rms) width, w , of the cloud in both the $\hat{y}$ and $\hat{z}$ directions. The extracted widths are binned by time of flight, and the error bars are calculated from the standard error about the mean. We fit the data to $w(t)^2 = w_0(t)^2 + v_{rms}^2 t^2$, where w_0 is the initial rms width of the cloud, $v_{rms} = \sqrt{k_B T / m}$ is the rms velocity, k_B is Boltzmann's constant, m is the atomic mass, and T is the temperature of the atomic cloud. The temperature is 7.8(9) mK and 4.6(4) mK for $\hat{y}$ and $\hat{z}$, respectively, where the errors in parentheses are one standard deviation, which is consistent with a conventional six-beam MOT [118]. The temperature is not equal in the two dimensions because the diffusion coefficient and velocity damping constant are different in the axial and radial directions in a grating MOT. For our MOT, the ratio of the temperatures along $\hat{y}$ and $\hat{z}$ is 1.7(2), consistent with the ratio of 1.9 from the theory in Ref. [49].

While the discussion has focused on trapping ^{88}Sr , strontium has a number of stable isotopes. The isotope abundances for strontium are 82.58 %, 7.00 %, 9.86 %, and 0.56 % for ^{88}Sr , ^{87}Sr , ^{86}Sr , and ^{84}Sr , respectively. Our setup can also trap around 7×10^5 atoms of ^{86}Sr at a source current of 13 A, consistent with the abundances above. Likewise, we would also expect to trap around 5×10^5 atoms of ^{87}Sr , but were unable to realize a MOT of ^{87}Sr . The hyperfine structure of ^{87}Sr poses at least two complications. First, we might not have sufficient repump power to adequately address all necessary hyperfine transitions [8]. Second, the hyperfine structure combined with the non-trivial geometry and polarizations of the grating MOT may significantly weaken the already limited transverse confining

forces [44, 113]. The theoretical details of the latter are beyond the scope of this work and will be presented in a future publication.

5.2.5 Discussion

We have realized a grating MOT of alkaline-earth atoms in a compact 2.4 L apparatus. Our permanent magnet design supplies the necessary field gradients for the MOT and allows for a degree of tunability, while the differential pumping tube limits outgassing from the hot source. The MOT traps up to 5×10^6 atoms of ^{88}Sr at a loading rate of $4 \times 10^6 \text{ s}^{-1}$, with a lifetime of approximately 1 s. This performance is comparable to vapor loaded, six-beam strontium MOTs [118, 125]. We also observe MOTs of ^{86}Sr with 7×10^5 atoms, consistent with the relative isotopic abundance.

In the future, upgrades to our apparatus could be made to improve the performance of the MOT and decrease the size of the system. We could improve the quality of the vacuum and reduce the size of the apparatus by replacing the ion pump with a hybrid NEG/ion pump. Improving the quality of the vacuum would increase atom number and lifetime, potentially allowing us to observe a MOT of ^{84}Sr . The system could be further miniaturized by using a fiber-coupled and photonically integrated chip to expand the MOT beam to the appropriate size without additional optics [126, 127].

With additional upgrades to our apparatus, we would be able to transfer atoms to a second stage MOT operating on the narrow $^1\text{S}_0 \rightarrow ^3\text{P}_1$ transition at 689 nm. The grating has good diffraction efficiency at both the 461 nm and 689 nm cool-

ing wavelengths. In addition, 43% of the capture volume created by the 689 nm diffracted beams overlaps with the 461 nm capture volume, facilitating transfer between the MOTs. As discussed in Ref. [113] and Ref. [128], the diffracted beams of a grating MOT have a complicated polarization projection onto the $\hat{z}$ axis. The mixed polarization projection of the diffracted beams reduces the confining force, which might limit the second stage MOT because the force due to gravity is no longer negligible. The second stage MOT also requires a low magnetic field gradient to operate, which we could achieve by incorporating electromagnets with our permanent magnet assembly. Due to its low field gradient, the second stage MOT is highly sensitive to stray magnetic fields, thus we would also need to incorporate shim coils to ensure proper positioning. The electromagnets and shim coils would also allow us to null the magnetic field to allow operation of optical clocks. We plan to make these upgrades to the apparatus and attempt second stage cooling in future experiments.

The implementation of field-deployable quantum devices relies on compact systems. Alkaline-earth-based quantum sensors have been proposed as platforms for atom interferometers and atomic clocks. Compact interferometers could be used for inertial navigation [32], and gravitational wave detection in space [34]. Deployable networks of optical clocks will be important for improved time and frequency metrology [36], and tests of fundamental physics [35]. Our results show that alkaline-earth grating MOTs are a promising step towards the development of compact optical clocks and other quantum devices.

5.2.6 Acknowledgements

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5.2.7 Data Availability

The data that support the findings of this study are available from the corresponding author upon reasonable request.

Chapter 6: Future Ideas and Conclusion

Our new strontium apparatus opens the way for more experiments in the lab. Our new high flux source provides high atom numbers, our chamber has a large amount of optical access for additional beams, and our new imaging system allows for the creation of high resolution optical potentials, and imaging. With these improvements, the lab is ready to move forward with the next set of experiments. We have a variety of ideas for future directions to take the laboratory, including quantum simulation in a 1D box trap with $SU(N) \times S_n$ symmetry, Rydberg spin squeezing of the optical clock transition, the creation of sub-wavelength attractive potentials with three-level systems, and a variety of new direction for the grating MOT experiment. These possible experiments are outlined in this chapter.

6.1 Box Trap

Quantum simulation relies on the creation of hamiltonians that match that of a system to be studied. Hamiltonians with $SU(N) \times S_n$ symmetry can be used to study spin models with infinite range interactions [16] or estimate the spectrum of an N dimensional density matrix that is encoded in the nuclear spin of each atom in the box trap [17]. As mentioned in the Sec. 1.1, the nuclear spin of ^{87}Sr creates an

$SU(N)$ symmetry, where $N \leq 10$, in the ground and clock states. This symmetry can be expanded to $SU(N) \times S_n$, where n is the number of atoms, by putting the atoms in a 1D box trap potential.

We would create this system by loading ultracold atoms into a tight dipole trap, such that the atoms only occupy the lowest energy level of the harmonic oscillator potential in the transverse x and y directions. In the z direction we would add the walls of the 1D box trap with a repulsive dipole trap, and allow the atoms to occupy any motional level of the box trap. By initially spin polarizing the gas, we could ensure that no level is doubly occupied. The wavefunctions for the levels of the box trap are

$$\psi_j(z) = \sqrt{2/L} \sin(j\pi z/L)$$

We can model the interactions between atoms in the box trap with a pseudo-potential $V(\vec{r}) \propto \delta(\vec{r})$, where $\vec{r}$ is the separation between the atoms. Then the interaction energy between atoms initially in states $\psi_i(z)\psi_j(z')$ scattering to states $\psi_{i'}(z)\psi_{j'}(z')$ becomes

$$U \propto \int_0^L \psi_{i'}(z)\psi_{j'}(z)\psi_i(z)\psi_j(z)dz$$

where we have integrated over the transverse directions and integrated the delta function over z' . From energy conservation considerations and this integral, the result is that U is a constant when $(i, j) = (i', j')$ or (j', i') , and zero otherwise. Thus this interaction can swap states of atoms between two box trap levels, and the

energy is independent of the levels involved, this is associated with a permutation symmetry S_n for the n levels of the box trap that are occupied. For the ground and clock state where the nuclear spin is decoupled from the electronic angular momentum, this implements an $SU(N) \times S_n$ symmetry.

We have started to implement the box trap in the laboratory. Initially, we planned on using a digital mirror display (DMD) to generate the walls of the box trap, but the low efficiency at 461 nm prompted us to switch to a printed mask. We have gotten the mask printed with a variety of sizes to generate an $\sim 10 \mu\text{m}$ length trap with our current imaging system. We plan on using the old 461 nm system to generate the blue-detuned repulsive dipole trap that will create the walls of the trap, and using the high-resolution imaging system to image this potential onto the atoms. For these experiments to work, the walls of the box trap must be high enough, and the atoms cold enough, that the occupied box trap levels maintain the S_n symmetry. In addition, the transverse dipole trap must be deep enough, that at the given temperature only the lowest motional level of this trap is excited.

6.2 Rydberg Spin Squeezing

One of the limits on strontium optical clocks is the quantum projection noise related to measuring the clock transition on independent atoms [129]. To beat this limit and decrease the noise on the measurement entanglement has to be added to the system so the measurements are no longer independent. By adding interactions, and generating “spin-squeezing” in the ensemble of atoms, the noise in the measurement

of the proportion of atoms in the ground $|g\rangle$ or clock $|c\rangle$ states is improved at the expense of the noise in the phase (when operating on the equator of the Bloch sphere).

One way to add interactions and implement this spin squeezing is by dressing the clock state with a Rydberg state $|r\rangle$ so that the Rydberg dressed clock state is $|c'\rangle \sim |c\rangle + \epsilon|r\rangle$ [28, 29]. Due to the large electric dipole moments of Rydberg atoms, they interact strongly through the Van der Waals force. This interaction can be used to implement entanglement and spin squeezing in an ensemble of atoms. Our laboratory has a laser operating at ~ 317 nm, mentioned in Sec. 3.3.6 to address Rydberg levels off the clock transition and implement this dressing for investigations of spin squeezing in strontium optical clocks.

6.3 Sub-wavelength lattices

Most often potentials for quantum simulation in cold atom systems are created with optical lattices and dipole traps. In general, the spatial scale of these potentials is limited by the wavelength of the light used. As an alternative, the Rb/Yb mixtures group at the JQI has used a three-state system to create effective potentials with structure smaller than the wavelength of the coupling lasers [130]. The basic idea of this scheme uses a homogeneous probe field Ω_p that couples the ground state to the intermediate state, and a spatially varying control field Ω_c that couples the intermediate state to the third state. This system then has a dark state that varies in space, when $\Omega_p \gg \Omega_c$ the dark state is the ground state, and when $\Omega_c \gg \Omega_p$ the

dark state is the third state. In between these two extremes the dark state (and energy) varies quickly in space creating the sub-wavelength structure of the lattice. So far this scheme has only been used to create repulsive sub-wavelength structure, but it can create attractive potentials with the right parameters [30]. Based on my discussions with Sarthak Subhankar, it seems that attractive potentials might be easier to achieve with strontium than with ytterbium, motivating this as a potential experiment for the new strontium laboratory. Sub-wavelength attractive potentials have a variety of applications from expanding the parameter space of quantum simulators, to improving quantum information processing.

6.4 Grating MOT Ideas

Our strontium grating MOT started as a side project, but it has proven to be an interesting avenue to explore. There are several improvements and new directions the project could take. The most immediate goal is to find a red MOT of ^{87}Sr in this system. As mentioned in Sec. 2.3, the interaction of the narrow linewidth, hyperfine structure, and complex polarizations in this system make it hard to intuitively understand which conditions will provide stable operation. Despite this, our simulations have indicated that it should be possible to operate a red grating MOT of ^{87}Sr by stirring on the $F = 9/2 \rightarrow 7/2$ hyperfine transition. So far we have not found a signal, but are hopeful that by increasing the power, narrowing the parameter space to search, and using the higher signal to noise ratio of fluorescence imaging we will find something soon.

The next step to further prove the usefulness of our strontium grating MOT system, could be to build an optical clock or atom interferometer. These systems would require a redesign of the chamber to allow for more optical access and the inclusion of an optical lattice or in the case of the interferometer a tower to increase the path length. There are also upgrades that could be made to the apparatus regardless of intent including improving the size and efficiency of the pumping, re-implementing an atomic shutter that will not jam, and redesigning the coils and coil electronics to produce both the red and blue MOT gradients with a single set of coils.

Another direction could be combining the strontium grating MOT and lithium grating MOT from the CAVS laboratory at NIST. The lithium grating MOT is used as a pressure standard by looking at atom loss out of a magnetic trap, with the ability to measure pressure more accurately than ion gauges [131]. By varying the background pressure in the chamber, and measuring it with the lithium trap as a pressure standard, we could measure the loss of strontium atoms and extract the scattering cross section. Background gas collisions are estimated to cause shifts to optical clocks at the 10^{-19} level [12, 13], so by doing this measurement we would be showing a proof of principle for the CAVS system, and aid the next generation of optical clocks.

6.5 Conclusion

During my time in graduate school, I was fortunate to help complete an experiment on a mature apparatus. We precisely measured the isotope shifts of the clock and intercombination transitions in strontium, finding a deviation from King linearity for ^{87}Sr . Additionally, we were able to build a new and improved strontium laboratory for a new generation of experiments. We were able to work through all the complexities of laser cooling strontium, ending with ultracold strontium atoms loaded in optical dipole traps, and set up the apparatus for the future of the group. While building up this apparatus, we were lucky to have Dan Barker suggest we make a strontium grating MOT. We developed the first alkaline earth grating MOT, a step towards realizing practical compact quantum devices that take advantage of the unique properties of strontium. This project gave us a lot to work on while the COVID-19 pandemic slowed down the speed of progress on the main apparatus. From our continued work on this system, we were able to complete the transfer between the broad and narrow line grating MOTs and show that SWAP works efficiently in the complex polarizations of a grating MOT.

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