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Quantum State Preparation in an Optical Lattice

by

Steven Eugene Hamann

A Dissertation Submitted to the Faculty of the
DEPARTMENT OF PHYSICS
In Partial Fulfillment of the Requirements
For the Degree of
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THE UNIVERSITY OF ARIZONA

1998
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DEDICATION

This work is dedicated to my parents for their many years of patience.
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This dissertation reports on quantum state preparation of cesium atoms in a two-dimensional optical lattice, by resolved-sideband Raman cooling. An optical lattice is a periodic potential produced by the light shift interaction between an atom and light field. Laser cooled atoms can become strongly localized about the bottom of potential wells in an optical lattice, where they occupy a discrete spectrum of bound vibrational energy levels. The distribution over vibrational levels of atoms in the lattice is characterized by the mean vibrational excitation, $\bar{n}$. In an optical lattice, absorption and emission of photons from lattice beams causes $\bar{n}$ to increase in time. This source of heating is always present, but its rate can be greatly reduced in a lattice detuned far from the atomic resonance. Sideband cooling is an efficient means of transferring atoms from higher into lower-lying vibrational levels and, thus, it reduces $\bar{n}$ for the ensemble. If the sideband cooling rate is much greater than the heating rate, then $\bar{n}$ approaches zero and virtually all atoms are in the lowest vibrational level in their potential wells. Our sideband cooling scheme involves stimulated Raman transitions between bound states in the potential wells of a pair of magnetic sublevels, followed by optical pumping, for a net loss of one quantum of vibration per cooling cycle. The process accumulates 98% of atoms in the ground vibrational level of a potential well associated with a single Zeeman substate. Each atom in the lattice is then very close to a pure state. For two-dimensional lattice with sideband cooling we find $\bar{n}_x \approx \bar{n}_y \approx 0.008(16)$.

Various issues related to state preparation and sideband cooling are also discussed in the context of a one dimensional lin $\perp$ lin optical lattice. These include improvement of laser cooling in a near resonance lattice by application of weak magnetic fields, transfer of atoms from near into far off-resonance lattices, and heating rates in far off-resonance lattices.
CHAPTER 1
INTRODUCTION

It has long been known that light can be used to manipulate atomic motion. With the introduction of lasers, it became practical to cool and confine atoms with laser light [1, 2]. First proposals for cooling and trapping were based on radiation pressure and were expected to produce temperatures no less than \( T_D = \hbar \Gamma / k_B \), the so called Doppler limit, where \( \Gamma \) is the linewidth of the atomic transition excited by the light field. New interest in laser cooling was generated by the surprising discovery of temperatures below the Doppler limit [3], and a theory of polarization gradient laser cooling was proposed [4]. In this theory, cooling is ascribed to a delicate interplay between optical pumping, which causes transitions among internal states of the atom, and the motion of the atom through periodic potentials generated by its dipole interaction with the light field. Due to its periodic nature, the spatially varying light shift experienced by an atom in a light field, which is generated by the interference of laser beams, is referred to as an optical lattice. The potential wells in a lattice are typically so shallow that they support only a few bound states for an atom. In a lattice made of near resonance laser light, cooling is so efficient that an atom can become confined in an individual potential well, where its mean energy is comparable to the spacing between bound levels; it must then be treated quantum mechanically [5, 6, 7]. In a lattice detuned far from the atomic transition, dissipation due to photon absorption and emission is suppressed, and the system is well suited to experiments involving quantum state preparation and control.

The objective of this dissertation is to describe experiments involving quantum state preparation of atoms in an optical lattice, most notably, resolved-sideband Raman cooling (sideband cooling). Sideband cooling is the central experiment in my Ph.D. work and the one for which I had primary responsibility.
The dissertation is organized as follows. Chapter two presents theoretical aspects of optical lattices and state preparation. It is by no means a general review of laser cooling theory, but a synopsis of concepts vital to understanding the experiments described later. Chapter three outlines the experimental apparatus and diagnostic techniques. It contains details which are not absolutely necessary for understanding the experiments discussed later. However, a major part of my contribution was construction of the apparatus which occupied about one year, and, so, it is dealt with at some length. Chapter four presents results from experiments in a one dimensional lin \( \perp \) lin optical lattice, which laid the groundwork for sideband cooling. These experiments include transferring laser cooled atoms from a near resonance into a far off-resonance optical lattice, determining heating rates in a far off-resonance lattice and improving the preparation of atoms in a near resonance optical lattice by application of magnetic fields. A description of the sideband cooling experiment and results are given in chapter five. The experiment is carried out in a two-dimensional lattice and some technical issues surrounding its construction and use are considered. Sideband cooling prepares atoms in the ground vibrational level of lattice potential wells, in pure quantum states. It is a good starting point for undertaking a variety of other experiments. Results for two such experiments are presented. The first involves squeezing of the momentum distribution of trapped atoms, and the second involves adiabatic cooling of sideband cooled atoms to achieve subrecoil temperatures. Finally, a brief summary of major results from the experiments described in chapters four and five is presented in chapter six.
CHAPTER 2
THEORETICAL BACKGROUND

2.1 Optical Lattices

2.1.1 Introduction

The light shift or ac Stark shift is a conservative interaction between an atom and a light field. It may be thought of as the polarization energy of the induced atomic dipole moment of an atom in a time varying electric field, the laser field [8]. It is known from atomic physics that an atom in a static electric field experiences a shift (called the dc Stark shift) of its internal energy levels.

The light shift energy of an atom depends upon the light intensity, the local polarization at the position of the atom and the atomic internal state. If the intensity or polarization of the field varies spatially in a periodic fashion, this gives rise to a periodic optical potential which will, in general, be different for each internal sublevel of the atom [9, 10]. It is possible to produce a periodically varying intensity or polarization by interfering laser beams. An optical lattice, then, is an interference pattern between laser beams in which an atom experiences a potential via the light shift. The primary force involved with Doppler cooling and magneto-optic traps is the spontaneous force which comes about due to photon scattering [1]. Photon scattering occurs in an optical lattice as well, but in a lattice the conservative part of the interaction between the atom and light is especially important. In a lattice formed by light detuned very far from the atomic resonance, photon scattering is greatly suppressed and it is then reasonable to attempt experiments which involve preparation or coherent evolution of the atomic internal state or center-of-mass motion such as resolved-sideband Raman cooling or momentum squeezing.
2.1.2 Sisyphus Cooling

In the presence of near resonance lattice light an atom undergoes optical pumping between internal magnetic sublevels. The periodic potentials of an optical lattice together with optical pumping provide for the possibility of cooling well below the Doppler limit. The mechanism behind this called, Sisyphus cooling, was described in the 1D lin ⊥ lin optical lattice for a $J_g = 1/2$, $J_e = 3/2$ atom [4]. Sisyphus cooling relies upon transitions between magnetic sublevels of the ground state, caused by optical pumping, as the atom moves along the potentials in the lattice. This cooling removes kinetic energy from the atoms so efficiently that eventually they become deeply trapped in the lattice potential wells where their center of mass motion must be treated quantum mechanically. Then the atoms occupy a discrete spectrum of bound vibrational energy levels. In a potential well with many bound levels, those with lowest energies, much less than the peak-peak potential depth of the well, are said to be "low lying". The lower the energy level occupied by an atom, the smaller its RMS position spread about the well bottom and the more tightly bound it is said to be.

2.1.3 Cesium

For all our experiments we use Cesium which is a group 1A alkali atom with a single valence electron. An alkali has the character of a one-electron atom with a relatively simple sublevel structure, which is an advantage for laser cooling. In its ground state the valence electron of Cesium is in the $6S_{1/2}$ level. We laser cool between that and the excited $6P_{3/2}$ state, which is the $D_2$ line of Cesium with a wavelength of $\lambda = 852.1$ nm. Coupling between the spin-orbit angular momentum and that of the nucleus divides these levels into nondegenerate hyperfine levels (figure 2.1). There are four, $F' = 2, 3, 4$ and 5 in the excited $6P_{3/2}$ state, and two, $F = 3$ and 4 in the $6S_{1/2}$ ground state. The spacing between the ground state hyperfine levels is about 10 GHz while spacings among excited state hyperfine levels are hundreds of MHz. For Sisyphus cooling in a near resonance optical lattice the laser is detuned red of the $F = 4$ to $F' = 5$ transition, referred to as the "cooling transition". The natural
Figure 2.1: Cesium energy level structure for the $D_2$ line. Excited and ground state hyperfine levels are labeled by $F'$ and $F$, respectively. Zeeman sublevels with $m \geq 0$ corresponding to some hyperfine levels are shown. Dashed lines indicate the detuning of the NRL and FRL from the $F = 4$ to $F' = 5$ cooling transition.

Linewidth of the transition is denoted by $\Gamma$ and has the value $2\pi \times 5.22$ MHz. The hyperfine levels are further divided into Zeeman sublevels labeled by $m$ which run from $m = -F$ to $+F$ in integer steps. These Zeeman or magnetic sublevels are degenerate in the absence of magnetic fields. Though they add complexity to the optical lattice, transitions among these sublevels allow for Sisyphus cooling and, as we shall see, for resolved-sideband Raman cooling as well.
2.1.4 The Potential Operator

The potential operator for an atom in a low intensity laser field in the presence of an external uniform constant magnetic field is given by [11],

$$\hat{U}(\mathbf{x}) = -\mathbf{E}_L^*(\mathbf{x}) \cdot \hat{\alpha} \cdot \mathbf{E}_L(\mathbf{x}) - \hat{\mu} \cdot \mathbf{B}$$

(2.1)

where $\mathbf{E}_L(\mathbf{x})$ is the spatial part of the laser field, $\hat{\alpha}$ is the atomic polarizability tensor, $\hat{\mu}$ is the magnetic dipole operator and $\mathbf{B}$ is the magnetic field. The magnetic dipole operator has the form,

$$\hat{\mu} = -\gamma g \hat{F},$$

(2.2)

where $\gamma g$ is the gyromagnetic ratio, and $\hat{F}$ is the total angular momentum operator.

The atomic polarizability tensor in the far off-resonance limit is given by,

$$\hat{\alpha} = -\sum_{F'} \frac{\hat{d}_{FF'} \hat{d}_{F'F}}{\hbar \Delta_{F,F'}},$$

(2.3)

where $\hat{d}_{FF'}$ is the electric dipole operator between the ground and excited hyperfine states labeled by $F$ and $F'$, respectively. For our case $F = 4$ since we laser cool atoms in this hyperfine ground state. Further, $\Delta_{F,F'}$ is the detuning of the laser from the $F$ to $F'$ transition and the summation is over the three excited states $F' = 3, 4$ and 5. Each excited state hyperfine level is divided into $2F' + 1$ Zeeman sublevels. A convenient basis to utilize is the basis of eigenstates of $\hat{F}^2$ and $\hat{F}_z$ where the quantization axis of the lattice is always taken along the $z$ axis. These states will be denoted by $|F,m\rangle$ where $F(m)$ labels the eigenstate associated with the operator $\hat{F}^2(\hat{F}_z)$. The electric dipole operator, $\hat{d}_{FF'}$, implicitly contains a summation over all allowed transitions between ground and excited state Zeeman sublevels. In general, the allowed transitions depend upon the polarization present in the light field. In the basis of spherical polarization unit vectors $\hat{e}_q$, and far off-resonance ($\Delta_{F,F'} \gg \Gamma$), the components of the polarizability tensor taken between unit vectors $\hat{e}_q$ and $\hat{e}_{q'}$ can be expressed as,
\[ \hat{\alpha}_{q,q'} = \tilde{\alpha} \sum_{F',m} \frac{\Delta F_{\text{max}},F'_{\text{max}}}{\Delta F,F'} f_{F,F'} \cdot c_{F',m+q}^{F',m+q'} c_{F,m}^{F,m+q} \times |F, m + q - q' \rangle \langle F, m|, \]

where \( \tilde{\alpha} \) is the characteristic polarizability scalar of the transition \(|J = 1/2\) to \(|J' = 3/2\) in this case, \(f_{F,F'}\) is the relative oscillator strength of the decay from \(|F'\rangle\) to \(|F\rangle\), \(c_{F',m}^{F',m'}\) is the Clebsch-Gordan coefficient for the transition \(|F, m\rangle\) to \(|F', m'\rangle\) and \(F_{\text{max}}\) labels the hyperfine level in the manifold \(\{F\}\) with largest eigenvalue. For our case \(F_{\text{max}} = 4\) and \(F'_{\text{max}} = 5\). In the spherical basis \(q = \pm 1\) or 0, which refers to \(\sigma^\pm\) and \(\pi\) polarized light, respectively. The matrix elements of the light shift operator are given by \(E_L^*(x) \cdot \hat{\alpha} \cdot E_L(x)\) taken between ground state eigenstates \(|F, m_i\rangle\) and \(|F, m_j\rangle\). Inspection of equation 2.4 and consideration of every possible combination of polarizations (or values of \(q\) and \(q'\)) reveals the only nonvanishing matrix elements to be those which connect states \(|F, m_i\rangle\) and \(|F, m_j\rangle\) where \(\Delta m \equiv (m_i - m_j) = 0, \pm 1\) or \(\pm 2\). Terms for which \(\Delta m = 0\) lie along the diagonal of the light shift operator. Off-diagonal terms that couple state for which \(\Delta m = \pm 1\) or \(\pm 2\) result from stimulated Raman transitions between ground state magnetic sublevels induced by appropriate combinations of \(\sigma\) and \(\pi\) polarized light. Since the polarization can vary greatly in the lattice on distance scales of order \(\lambda\), these matrix elements which depend upon local polarization will also vary on such scales. So the strength of the coupling between different Zeeman sublevels due to Raman transitions can be strongly connected to the position of an atom in the lattice.

2.1.5 1D Lin∥Lin Optical Lattice

At this point it is useful to specialize to the particular case of a 1D Lin∥Lin optical lattice, both because it is a relatively simple system which offers insight into general features of optical lattices, and because several experiments to be discussed were carried out in this geometry. The 1D Lin∥Lin optical lattice is formed by two counter-propagating beams of equal intensity with orthogonal linear polarizations.
Figure 2.2: Gradient of ellipticity for a 1D lin ⊥ lin optical lattice aligned along the z axis. Unit vectors $e_+$ and $e_-$ correspond to pure $\sigma^+$ and $\sigma^-$ polarizations, respectively, while $e_\perp = (1/\sqrt{2})(\hat{x} + \hat{y})$ and $e_\parallel = (1/\sqrt{2})(\hat{x} - \hat{y})$.

This geometry gives rise to a gradient of ellipticity which consists of two standing waves of $\sigma^+$ and $\sigma^-$ light spatially offset by $\lambda/4$. The local polarization varies from pure $\sigma^+$ to pure $\sigma^-$ in a quarter wavelength while the intensity is uniform. For such a lattice oriented along the z axis and for a particular choice of relative phase between the counterpropagating beams, the spatial part of the field can be written as,

$$\mathbf{E}(z) = \sqrt{2}E_1\{\cos(k_Lz)\, e_+ - i\sin(k_Lz)\, e_-\}$$

where $E_1$ is the amplitude of a single beam and $k_L = 2\pi/\lambda$ is the laser wavenumber. Note that it is expressible in terms of polarization vectors $e_+$ and $e_-$ only, which is true in general for a 1D lattice. Thus, the matrix elements of the light shift operator are nonzero only for diagonal terms and for terms coupling $|F,m_i\rangle$ and $|F,m_j\rangle$ where $\Delta m = \pm 2$. The diagonal elements of the light shift operator are the nine diabatic potentials, $U_{F,m}(z)$, associated with the nine eigenstates $|F,m\rangle$ and are given by [12],

$$U_{F,m}(z) = \frac{\hbar \Gamma}{8E_R} \frac{I}{I_0} \sum_{F',m'} \frac{\Gamma}{\Delta_{F,F'}} f_{F,m}^{F',m'} [(m' - m)\cos(2k_Lz) + 1]$$

in the far off-resonance limit (figure 2.3). Here $I$ is the intensity of one beam, $I_0$ is the saturation intensity of cesium, $E_R = \hbar^2 k_L^2/2M$ is the recoil energy in terms of the...
atomic mass, $M$, and $f_{F'm'}$ is the oscillator strength for the transition from the state $|F, m\rangle$ to the state $|F', m'\rangle$. Sublevels $|F, m \geq 1\rangle$ have minima at regions of pure $\sigma^+$ light while $|F, m \leq -1\rangle$ have minima at regions of pure $\sigma^-$. A useful quantity to define is the single beam light shift, $U_1$, which is the light shift experienced by an atom with a Clebsch-Gordan coefficient of one, in a single lattice beam. The lattice beams are taken to be linearly polarized running waves for this definition.

Adiabatic potentials which take into account Raman coupling are found by diagonalizing the light shift operator at each position, $z$, in the lattice [13](figure 2.4). This gives rise to new adiabatic eigenstates which can be expressed as linear superpositions of the basis functions $|F, m\rangle$. Therefore, an adiabatic potential cannot be associated with a single state $|F, m\rangle$. However, the eigenstate associated with the lowest lying adiabatic potential has almost entirely $|4, m = +4\rangle$ character in regions of the lattice around positions of pure $\sigma^+$ light. In such locations the adiabatic potential is nearly identical to the diabatic potential associated with the magnetic sublevel $|F, m = +4\rangle$, denoted by $U_{4,+4}(z)$. As the atom moves away from a region of pure $\sigma^+$ light, other magnetic sublevels begin to be strongly added into the adiabatic eigenstate. This creates an avoided crossing midway between regions of pure $\sigma^+$ and $\sigma^-$ light. Finally, in the region around pure $\sigma^-$ light the eigenstate and potential are nearly identical to $|F, m = -4\rangle$ and $U_{4,-4}(z)$, respectively. In general, if Raman coupling is significant, the adiabatic potentials best describe the lattice, as long as the atoms move slowly enough through the avoided crossing regions to undergo Raman transitions. In any case, for an atom tightly bound around regions of pure $\sigma^+$ or $\sigma^-$ light, the adiabatic and diabatic potentials and eigenstates are nearly identical.

2.1.6 Far Off-Resonance Optical lattices

Inspection of equation 2.6 reveals the dependence of the optical potentials on the intensity and detuning of the laser field. In the very far off-resonance limit where the lattice detuning is much greater than the frequency spacing between excited state hyperfine levels, the optical potential depth can be characterized by a single
Figure 2.3: Diabatic Potentials for the even (solid lines) and odd (dashed lines) magnetic sublevels in a 1D lin ⊥ lin optical lattice for lattice parameters (a) detuning, $\Delta = -20\Gamma$ and single beam light shift, $U_1 = 90E_R$ and (b) $\Delta = -3831\Gamma$ and $U_1 = 75E_R$. 
Figure 2.4: Adiabatic Potentials in a 1D lin ⊥ lin optical lattice for detuning, \( \Delta = -20\Gamma \) and single beam light shift \( U_1 = 90E_R \).

detuning, \( \Delta_{F_{\text{max}},F'_{\text{max}}} \). In that case,

\[
U_{pp}^{m} \propto \frac{I}{\Delta},
\]  

(2.7)

where \( U_{pp}^{m} \) is the peak-peak diabatic potential depth for an atom in the \( m^{th} \) magnetic sublevel of the \( F = 4 \) ground state, and the cumbersome subscript \( F_{\text{max}},F'_{\text{max}} \) has been dropped. In this same limit the rate at which atoms scatter light from the lattice beams, \( \gamma_s \) is given in terms of \( U_{pp}^{m} \) by,

\[
\frac{\hbar \gamma_s}{E_R} = \frac{3}{2} \left( \frac{U_{pp}^{m}}{E_R} \right) \frac{\Gamma}{|\Delta|}.
\]  

(2.8)

It is evident that if \( I \) and \( \Delta \) are proportionally increased, the peak-peak potential depth is unchanged while the scattering rate decreases. Scattering of lattice photons, while necessary for Sisyphus cooling, is problematic for experiments such as momentum squeezing or sideband cooling. When an atom spontaneously emits a photon,
any coherent motion it is undergoing is interrupted. Therefore, only coherent processes which occur on time scales much shorter than $1/\gamma_s$ can be detected. Photon scattering also imparts random momenta to bound atoms which heats them into higher lying vibrational levels, eventually permitting escape from the well. If preparation of atoms in low lying vibrational levels is a goal, this heating is a competing process and the preparation can only be accomplished if the cooling rate is much greater than the heating rate which depends upon $\gamma_s$. The motivation for working in far off-resonance lattices then is clear. In such lattices the photon scattering rate is suppressed and the time scale $1/\gamma_s$ is increased. In general, if Sisyphus cooling is not needed and a low photon scattering rate is required, it is best to detune the lattice as far as available laser power permits.

2.2 Magnetic Fields in an Optical Lattice

The potential operator for an atom in an external magnetic field is given by the second term on the right hand side of equation 2.1. Consider the longitudinal component of this operator along the quantization axis $z$, given by $(\gamma g B_z) \hat{F}_z$. This component is diagonal in the basis $\{|F, m\rangle\}$ and produces matrix elements of the form,

$$\langle F, m | \gamma g B_z \hat{F}_z | F, m \rangle = (g_F \mu_B m) B_z,$$

(2.9)

where $g_F$ and $\mu_B$ are the Lande g-factor and Bohr magneton, respectively. This represents an uniform energy shift of the diabatic potentials which differs for each sublevel $|F, m\rangle$ based upon $B_z$ and eigenvalue, $m$. For example, if $B_z < 0$, then the diabatic potential $U_{4,4}$ is uniformly shifted down in energy more than $U_{4,3}$. The effect upon the lowest lying adiabatic potential is similar. In a positive magnetic field the potential minimum associated with $|4, m = 4\rangle$ is shifted up in energy and that associated with $|4, m = -4\rangle$ is shifted down (figure 2.5). This causes the potential well associated with $|4, m = -4\rangle$ to become deeper with more bound states and that associated with $|4, m = 4\rangle$ to become shallower with fewer bound states. In a near resonance lattice the result of this is accumulation of atoms in the deeper well in
the stretched state $|4, m = -4\rangle$ [14, 15].

The transverse components of the magnetic potential operator along $x$ and $y$ have the form $(\frac{1}{2}B_x)(\hat{F}_+ + \hat{F}_-)$ and $(\frac{1}{2}B_y)(\hat{F}_+ - \hat{F}_-)$, respectively, where

$$\hat{F}_\pm |F, m\rangle = \hbar \sqrt{F(F+1) - m(m \pm 1)} |F, m \pm 1\rangle .$$

(2.10)

Operators $\hat{F}_+$ and $\hat{F}_-$ give rise to off-diagonal matrix elements that couple magnetic sublevels which differ by $\Delta m = \pm 1$. As we saw, such couplings are also produced by appropriate combinations of $\sigma$ and $\pi$ polarized light. However, in this case the couplings are not spatially varying since $B$ is uniform. In a near resonance lattice magnetic field induced coupling between ground state magnetic sublevels can increase the cooling efficiency and cause atoms to become more tightly bound in the adiabatic potential wells associated with the stretched states $|4, m = \pm 4\rangle$ [15]. Moreover, weak magnetic fields can increase the population of atoms in the lattice in a single stretched state and in lower lying vibrational levels. This represents
a limited type of state preparation, and is useful in experiments involving state selection or state preparation such as sideband cooling, which will be discussed.

2.3 Temperature

2.3.1 Kinetic Temperature

The kinetic temperature of a sample of atoms is related to the "spread" of momentum values present. One measure of the spread of a momentum distribution in one dimension is its variance given by

\[(\Delta p)^2 = \int_{-\infty}^{+\infty} p^2 \Pi(p) \, dp - \left( \int_{-\infty}^{+\infty} p \, \Pi(p) \, dp \right)^2,\]  

(2.11)

where \(p\) is the momentum and \(\Pi(p)\) is the normalized momentum probability distribution of the sample. If \(\Pi(p)\) is a Gaussian function,

\[\Pi(p) = \frac{1}{\sqrt{2\pi}\sigma_p} e^{-\frac{p^2}{2\sigma_p^2}},\]  

(2.12)

then, in one dimension, the temperature with respect to the momentum degree of freedom or kinetic temperature, \(T_K\), is defined in terms of the variance of this function, \(\sigma_p^2\), and is given by [16],

\[T_K = \frac{\sigma_p^2}{M k_B}.\]  

(2.13)

2.3.2 Distribution over Vibrational Levels

An atom tightly bound in a diabatic potential well of a 1D lattice can be approximated as 1D harmonic oscillator. The total eigenstate is then a product \(|F, m\rangle |n\rangle\), denoted \(|F, m, n\rangle\), where \(|n\rangle\) is the harmonic oscillator eigenstate of the \(n^{th}\) vibrational level. The vibrational temperature, \(T_{vib}\), of a sample of atoms in the lattice relates to their distribution over these vibrational levels. For a thermally excited system the vibrational temperature enters into the normalized occupation probability
function of the \( n^{th} \) energy level, \( \Pi_n \), in the following way [16],

\[
\Pi_n = \frac{e^{-\frac{E_n}{k_B T_{\text{vib}}}}}{\sum_m e^{-\frac{E_m}{k_B T_{\text{vib}}}}},
\]

(2.14)

where \( E_n \) is the energy of the \( n^{th} \) level. An important quantity related to \( T_{\text{vib}} \) called the mean vibrational excitation is defined as

\[
\bar{n} \equiv \sum_n n \Pi_n.
\]

(2.15)

It is an ensemble average of the occupation number \( n \). Note that if \( \bar{n} \) is known, \( T_{\text{vib}} \) can be determined from equation 2.14 and vice versa. The connection between \( T_K \) and \( T_{\text{vib}} \) can be found by considering the relationship between the momentum distributions of the vibrational levels and the measured momentum distribution of atoms released from the lattice. The lattice of bound atoms is a large ensemble of independent harmonic oscillators and so its measured momentum distribution is a population weighted sum of the momentum probability distributions of the individual vibrational levels. Therefore, the variance of momentum for all atoms in the lattice, \( (\Delta p)^2 \), is a population weighted sum of the variances of the momentum probability distributions of the vibrational levels, denoted \( (\Delta p_n)^2 \). This is expressed by

\[
(\Delta p)^2 = \sum_n \Pi_n (\Delta p_n)^2.
\]

(2.16)

The kinetic temperature is calculated from \( (\Delta p) \) by,

\[
\frac{(\Delta p)^2}{\hbar^2 k_B^2} = \frac{T_K}{T_R},
\]

(2.17)

where the recoil temperature, \( T_R = \hbar^2 k_B^2 / M k_B \), is introduced. The vibrational temperature enters through \( \Pi_n \). Since the form of \( \Pi_n \) is constrained by equation 2.14, then, as long as \( (\Delta p)^2 \) and \( (\Delta p_n)^2 \) are known quantities, the populations, \( \Pi_n \), can be calculated directly from equation 2.16 in terms of the parameter \( T_{\text{vib}} \). Then \( \bar{n} \), which is the quantity we generally use to describe the distribution over vibrational levels, can be calculated from equation 2.15.
For a harmonic oscillator equation 2.16 simplifies. In that case,

$$\langle \Delta p_n \rangle^2 = \langle n | \hat{p}^2 | n \rangle - \langle \langle n | \hat{p} | n \rangle \rangle^2 = \langle n | \hat{p}^2 | n \rangle ,$$

(2.18)

where $\hat{p}$ is the momentum operator and $|n\rangle$ is the $n^{th}$ harmonic oscillator eigenstate; $\langle n | \hat{p} | n \rangle = 0$. Defining $\langle n | \hat{p}^2 | n \rangle \equiv p_n^2$, then the expectation value of $\hat{p}^2$ in the $n^{th}$ eigenstate can be expressed in terms of its expectation value in the ground state,

$$p_n^2 = p_0^2(2n + 1),$$

(2.19)

where $p_0^2 = M\hbar\omega_{osc}/2$ is characterized in terms of the oscillation frequency. Plugging this into equation 2.16 and using the definition of $\bar{n}$ from equation 2.15 gives

$$\langle \Delta \hat{p} \rangle^2 = p_0^2(2\bar{n} + 1) \tag{2.20}$$

For a thermally excited harmonic oscillator the ensemble average momentum distribution is exactly Gaussian. If there are several bound states in the diabatic potential, it is nearly Gaussian in which case equation 2.13 can be invoked to express the kinetic temperature of the sample in terms of $\bar{n}$,

$$\frac{\langle \Delta \hat{p} \rangle^2}{p_0^2} = \frac{T_K}{T_0} = (2\bar{n} + 1).$$

(2.21)

The ground state kinetic temperature, $T_0 = p_0^2/Mk_B$ is introduced here. Note that if all atoms are in the ground state, $n = 0$, then $\bar{n} = 0$ and the measured temperature is the ground state kinetic temperature.

A final quantity of interest in the harmonic oscillator case is the Boltzmann factor, $q_B$, given by,

$$q_B \equiv e^{-\frac{\hbar\omega_{osc}}{k_B T_{vib}}} . \tag{2.22}$$

Using $q_B$ and the energy of the $n^{th}$ 1D harmonic oscillator level $E_n = (n+1/2)\hbar\omega_{osc}$, equation 2.14 can be written as,

$$\Pi_n = \frac{q_B^n}{\sum_m q_B^m} = (1 - q_B)q_B^n \tag{2.23}$$
Plugging this into equation 2.15 allows $\bar{n}$ to be expressed in terms of $q_B$ as

$$\bar{n} = (1 - q_B) \sum_n n q_B^n = \frac{q_B}{1 - q_B}$$  \hspace{1cm} (2.24)

In D dimensions equation 2.23 generalizes to

$$\Pi_n = (1 - q_B)^D q_B^n$$  \hspace{1cm} (2.25)

### 2.3.3 Harmonic and Anharmonic Approximations

In the harmonic approximation the oscillation frequency, $\omega_{osc}$, is determined by expanding the lattice potential to second order in $kz$. For a 1D lin $\perp$ lin optical lattice in the far off-resonance limit, the potential has the form,

$$U_m(z) = -U_{pp}^m \left[ 1 + \frac{1}{2} \cos(2kz) \right],$$  \hspace{1cm} (2.26)

where $U_{pp}^m$ is the peak-peak potential depth for an atom in the $m^{th}$ magnetic sublevel and $k = 2\pi/\lambda$. This can be expanded about $z = 0$ as

$$U_m(z) = -\frac{3}{2} U_{pp}^m + \frac{1}{2} (2U_{pp}^m k^2) z^2 - \frac{1}{3} U_{pp}^m k^4 z^4 + \frac{2}{45} U_{pp}^m k^6 z^6 + \cdots$$  \hspace{1cm} (2.27)

Keeping only the harmonic term we find the oscillation frequency in the well in terms of the peak-peak potential depth,

$$\frac{\hbar \omega_{osc}}{E_R} = 2 \sqrt{\frac{U_{pp}^m}{E_R}}.$$  \hspace{1cm} (2.28)

So the oscillation frequency depends upon the peak-peak potential depth of the particular well and this differs among Zeeman sublevels. To increase precision in the calculated vibrational temperature, in practice we make an anharmonic approximation. It is done by treating higher order terms in expansion 2.27 as a perturbation. This produces perturbed eigenfunctions which lead to adjusted values for the variances, $(\Delta p_n)^2$, and perturbed eigenvalues which alter the populations, $\Pi_n$. These adjusted variances and populations can then be used in equation 2.16 to calculate $T_{vib}$ and, therefore, $\bar{n}$ in the anharmonic approximation. We use the anharmonic approximation whenever calculating $\bar{n}$ from experimental data.
2.4 Resolved-Sideband Raman Cooling

2.4.1 Introduction

One type of state preparation in an optical lattice is to have all the independent atoms in identical quantum states, that is, in the same hyperfine and Zeeman sublevels and in the same vibrational level. In this case the lattice is an ensemble of atoms in separate or distinct but identical quantum states. Such a preparation may be achieved using resolved-sideband Raman cooling, which is a technique that has been applied in ion traps [17, 18], but which we specialize to far off-resonance optical lattices [11, 19]. The basic idea is this. Coupling is introduced which selectively transfers population from the \( n^{th} \) vibrational level of one stretched state, \(|F, m = 4, n\rangle\) into the \( n - 1 \) vibrational level of a different Zeeman sublevel, \(|F, m = 3 \text{ or } 2, n - 1\rangle\). Then optical pumping transfers the atom back into the stretched state without changing its vibrational level. So the two step sideband cooling process reduces the vibrational level of an atom by one unit, \( \Delta n = -1 \).

Harmonic oscillator eigenstates which differ by one unit have different parity and so the states \(|4, m = 4, n\rangle\) and \(|4, m = 3 \text{ or } 2, n - 1\rangle\) can only be coupled by matrix elements that have odd parity about their well bottoms which are aligned spatially. This rules out the use of uniform transverse magnetic fields, even though they can couple adjacent Zeeman sublevels for which \( \Delta m = \pm 1 \). As noted earlier, the couplings produced by a transverse field are uniform and so cannot connect states of different parity. We are left then with Raman coupling produced by appropriate combinations of \( \sigma \) and \( \pi \) polarized light which, as we saw, can give rise to off-diagonal matrix elements in the light shift operator which are spatially varying on the scale of \( \lambda \) and so have the potential for coupling states of different parity (figure 2.6). One possible method for inducing the desired population transfer with Raman coupling is to use two far off-resonance beams external to the lattice with appropriate polarizations which differ in frequency by an amount corresponding to the energy difference between states \(|4, m = 4, n\rangle\) and \(|4, m = 3 \text{ or } 2, n - 1\rangle\). The external beams couple these states via anti-Stokes stimulated Raman transitions (figure
However, it is simpler to capitalize upon the Raman coupling inherent to the lattice itself. The effect of this coupling can be greatly enhanced by application of a longitudinal magnetic field which Zeeman shifts the desired levels into energy degeneracy (figure 2.8)\cite{20}. This is the technique which we apply and have found to be successful. When sideband cooling, degeneracy between states $|4,m = 4, n \rangle$ and $|4,m = 3 \text{ or } 2, n - 1 \rangle$ is referred to as the first red sideband.

The sideband cooling cycle is completed by optically pumping the atom back into the stretched state, $|4,m = 4, n - 1 \rangle$. This is accomplished by application of a $\sigma^+$ polarized "pumper" beam which is resonant with the $F = 4$ to $F' = 4$ hyperfine transition (figure 2.9). An atom in the stretched state $|4, m = 4 \rangle$ cannot make a transition from $F = 4$ to $F' = 4$ due to the polarization of the beam. The transition rate from $F = 4$ to $F' = 5$ is negligible since the beam has low intensity and is detuned about 250MHz from the transition. Thus, an atom in the stretched state is dark to the pumper beam and will not heat up by scattering from it. Since decay to the $F = 3$ hyperfine ground state is possible when the atom excites to $F' = 4$, a
\( \sigma^+ \) polarized “repumper” beam resonant with the \( F = 3 \) to \( F' = 4 \) transition is also present. The repumper beam is much more intense than the pumper, but due to the large splitting between the hyperfine ground states, of order 10GHz, the \( F = 4 \) stretched state is dark to the repumper as well.

### 2.4.2 Cooling and Heating Rates

For sideband cooling to be effective the cooling rate must be much larger than the rate of increase of vibrational excitation, the heating rate, \( d\bar{n}/dt \). Since the pumper beam is on resonance, the optical pumping rate is high. So the rate of the two step cooling process is limited by the rate of Raman transitions involved. Thus, the Raman coupling, \( V_R \), must be strong enough that the Raman transition rate, \( |V_R|/\hbar \), satisfies the condition [11],
Figure 2.8: Zeeman shifting vibrational levels to the first red sideband by application of a longitudinal magnetic field. Vibrational levels of $|4, m = 3\rangle$ are shown as solid lines, while those of $|4, m = 4\rangle$ are dashed.

$$\left| V_R \right| \gg \frac{d\bar{n}}{dt}. \quad (2.29)$$

However, the Raman transition rate must not be so large that the vibrational energy levels are significantly broadened, since that would undermine the selective nature of the population transfer which is crucial to the sideband cooling process.

The heating rate is dominated by photon scattering. In the harmonic approximation and to first order in the square of the Lamb-Dicke parameter, $\eta = (k_Lz_0)$, it can be shown that [21],

$$\frac{d\bar{n}}{dt} \approx \gamma_s (k_Lz_0)^2, \quad (2.30)$$
Figure 2.9: Pumper and repumper beams with pure $\sigma^+$ polarization resonant with the $F = 4$ to $F' = 4$ and $F = 3$ to $F' = 4$ transitions, respectively.

up to a constant which is found by calculating the mean-squared momentum transfer per photon scattering event along the confining axes of the lattice. In this expression, $z_0$ is the RMS position spread of the harmonic oscillator ground state. Sideband cooling of tightly bound atoms can remove one vibrational quantum every few oscillation periods, $1/\omega_{osc}$. For effective cooling, then, the heating rate must be much less than the oscillation frequency. This is expressed by,

$$\gamma_s (k_L z_0)^2 = \gamma_s \left( \frac{E_R}{\hbar \omega_{osc}} \right) \ll \omega_{osc}$$

Using equations 2.8 and 2.28 for the case of a 1D lin $\perp$ lin optical lattice, the
inequality in equation 2.31 can be expressed as,

$$\frac{\Gamma}{|\Delta|} \ll \frac{8}{3}$$  \hspace{1cm} \text{(2.32)}$$

which is satisfied for large detunings.

2.4.3 Raman Coupling Strengths

As shown in the previous section, for sideband cooling to work the Raman coupling must be strong enough that the Raman transition rate satisfies equation 2.29. Consider Raman coupling between states \(|4, m = 4, n\rangle\) and \(|4, m = 2, n - 1\rangle\), denoted \(V_{2,4}\), in a 1D \(\text{lin} \perp \text{lin}\) optical lattice. In the far off-resonance limit it can be shown that,

$$V_{2,4} \propto \frac{\Gamma}{\Delta_{4,5}}$$  \hspace{1cm} \text{(2.33)}$$

where \(\Delta_{4,5}\) is the detuning of the lattice from the \(F = 4\) to \(F' = 5\) hyperfine transition. For a more detailed analysis see [11]. This behavior is due to a quantum destructive interference effect among the excited state hyperfine levels \(F' = 3, 4\) and 5 for transitions between ground states \(|4, m = 4\rangle\) and \(|4, m = 2\rangle\). Equation 2.33 indicates that Raman coupling between these states is weaker the further off-resonance the lattice is detuned, but that is exactly where sideband cooling should work best. So, prospects for sideband cooling are less favorable between these sublevels. No such destructive interference is present for Raman coupling between states \(|4, m = 4, n\rangle\) and \(|4, m = 3, n - 1\rangle\). For \(\Delta_{4,5} \rightarrow \infty\), Raman coupling between these states, \(V_{3,4}\), is independent of detuning which makes prospects for sideband cooling between these levels more favorable. Raman Coupling between states \(|4, m = 4\rangle\) and \(|4, m = 3\rangle\) requires a combination of \(\sigma\) and \(\pi\) polarized light. Since \(\pi\) polarized light cannot be introduced into a purely 1D lattice, sideband cooling between these levels is not possible in 1D.

In a 2D lattice it is possible, however, to introduce \(\pi\) polarized light. Consider the 2D lattice of (figure 2.10). It is formed by three coplanar beams of equal intensity with linear polarizations lying in the plane of the lattice and with 120 deg between
Figure 2.10: Two-Dimensional lattice geometry. The three lattice beams intersect at the origin with 120 deg between each and are polarized in the xy-plane.

each [22]. This geometry generates a lattice of 2D potential wells for the stretched state \(|4, m = 4\rangle\), which are nearly symmetric along the x and y axes (figure 2.11). Figure 2.12 shows the diabatic potentials associated with the states \(|4, m = 4\rangle\) and \(|4, m = 3\rangle\) along the y axis. A component of \(\pi\) polarization is introduced into the lattice by rotating the polarization of the beam propagating along the y-axis slightly out-of-plane. The electric field can then be written as,

\[
E_L(x) = \frac{E_1e^{-ik_Ly}}{\sqrt{2}}[-e_+\{1 + 2e^{iK_y}cos(K_x x)\} + e_-\{1 + 2e^{iK_y}cos(K_x x - 2\pi/3)\}] + e_0E_0e^{-ik_Ly}e^{i\phi},
\]

where \(E_1\) and \(E_0\) are the amplitudes of the in-plane and out-of-plane polarizations, and \(K_x = k_L \sin(\pi/3)\), and \(K_y = (1 + k_L \cos(\pi/3))\). Note that the \(\pi\)
Figure 2.11: Contour plot of two-dimensional diabatic potential associated with state $|4, m = 4\rangle$ for lattice detuning, $\Delta = -383 \Gamma$ and single beam light shift, $U_1 = 54 E_R$. Potential minima are shown in black.
polarization unit vector, \( e_0 \), is present. The strength of Raman coupling can be adjusted, in part, through the magnitude of the polarization component \( E_0 \). When this field is plugged into the potential operator (equation 2.1), the matrix elements for Raman coupling between \( |4, m = 4, n\rangle \) and \( |4, m = 3, n - 1\rangle \) are found to have even and odd contributions along \( x \) and \( y \). The odd contributions which couple the desired states are maximized for a relative phase shift of \( \phi = \pi/2 \) between the in-plane and out-of-plane polarization components [11]. Sideband cooling in a two-dimensional lattice is discussed further in chapter 5.
CHAPTER 3
EXPERIMENTAL APPARATUS

3.1 The Magneto-Optic Trap

3.1.1 Description

The potential wells of an optical lattice are extremely shallow compared to the kinetic energy of a typical room temperature atom. So to load an optical lattice with a large number of atoms requires a source of atoms which is substantially (about eight orders of magnitude) colder than room temperature. The magneto-optic trap (MOT) in combination with 3D optical molasses [1, 23, 2] is such a source, and is widely used for loading optical lattices or magnetic traps for example. The MOT is made of three pairs of counter-propagating beams of $\sigma+$ and $\sigma-$ light oriented along the $x$, $y$ and $z$ axes which all intersect at the origin. A pair of magnetic field producing “MOT coils” is centered about one axis such that the axis is normal to the plane of the coils and intersects the centers of the coils. The coils are also centered about the plane containing the alternate two axes. Current flows in the coils with the same magnitude but opposite sense of rotation (figure 3.1). The light is provided by a cooling laser detuned red of the $F = 4$ to $F' = 5$ hyperfine (cooling) transition. The MOT is set up in a vacuum chamber with a low pressure vapor of atoms at room temperature. Slow moving atoms, from the wings of the Gaussian distribution of velocity, which wander into the region of beam intersection are significantly Doppler cooled. The MOT coils produce a magnetic field which is zero at the origin and which increases in a positive sense along the $+x$, $+y$ and $+z$ directions and in a negative sense in the $-x$, $-y$ and $-z$ directions. This magnetic field Zeeman shifts the magnetic sub-levels of the atom, some closer to and some further from resonance with the cooling laser, based upon the sign and magnitude
of the magnetic field. These Zeeman shifts tune the scattering rate among the nine sub-levels and, in conjunction with the appropriate beam polarizations, produce a position dependent radiation pressure imbalance that always restores the atom to the zero of magnetic field. So making use of the internal structure of the atom, the MOT actually traps atoms at a particular point in space. For the trap to work in our case the atom must be in the $F = 4$ ground state. A beam is present to re-pump any atoms which decay to the $F = 3$ hyperfine ground state. This is referred to as the repumper beam.

Our MOT is built around an UHV chamber with a background vapor pressure of about $10^{-9}$Torr. The Cesium vapor pressure in the chamber is around $10^{-8}$Torr. Typical beam intensities, detunings and magnetic field gradients for loading the MOT are $6 I_0$ per beam, $-3.4 \, \Gamma$, and 10Gauss/cm, respectively. We trap a few
million atoms in a region of space roughly 300 µm in diameter, which corresponds to an atomic density of order $10^{11}$ atoms/cm$^3$.

### 3.1.2 3D Cooling and Transfer

The sequence of trapping, cooling and transfer of atoms to an optical lattice goes as follows. First, the MOT is loaded with a few million atoms in about three seconds. Then, it is compressed in size by increasing the magnetic field gradient to twice its initial value for 100 ms. Next, the magnetic field is switched off and the MOT beams are detuned further red to $-11.1 \Gamma$ and reduced in intensity by a factor of 10. This creates a three dimensional optical molasses which laser cools the atoms to around 3 µK in 10 ms. Finally, the atoms are transferred into a near resonance optical lattice by simply turning off the 3D molasses beams while, instantaneously turning on the superimposed optical lattice beams. At 3 µK the atoms are cold enough to efficiently Sisyphus cool in the near resonance lattice.

### 3.2 MOT Lasers

#### 3.2.1 External Cavity Lasers

High power, narrow line-width lasers are necessary for laser cooling. The line-width must be narrow enough that the various hyperfine transitions are well resolved. Although they do not offer the highest power available, diode lasers are an attractive option due to their relatively low cost and high reliability. Single mode diode lasers which can operate at up to half a watt output power are available, and when used in an external cavity configuration, they have sufficiently narrow line-width and are frequency tunable. Our external cavity tunable diode laser systems are built following a design by Wieman's group [24] which uses a diffraction grating in the Littrow configuration to provide feedback for the laser (figure 3.2). The $-1$ diffracted order is sent back into the laser and, so, the laser-to-grating distance defines an external cavity which narrows its line-width and imposes conditions on its frequency.

A piezoelectric transducer (PZT) attached to the grating allows for fine tuning of the external cavity length and, hence, laser frequency as well.
3.2.2 Frequency Stabilization of Lasers

In order to make a MOT or optical lattice it is obviously necessary to frequency stabilize the lasers relative to the laser cooling transition. One method of stabilizing lasers involves a technique called saturated absorption spectroscopy [25]. Consider figure 3.3. A beam passes through a cesium vapor cell, is retroreflected upon itself, passes through the cell again and impinges upon a photodetector circuit. As the atoms absorb light, the beam is attenuated and the change in the voltage output of the photodetector circuit is proportional to this attenuation. Since the beam is retroreflected, its left and right going parts are absorbed by atoms in the same spatial subset of the vapor cell. This arrangement leads to the photodetector signal shown in figure 3.4 as the laser frequency is scanned across the manifold of hyperfine transitions from the $F = 4$ ground state. The general properties of the photodetector signal can be explained as follows.

The linewdths of the transition and the laser are around 5 to 10 MHz. Let the laser frequency initially be $\nu_1$ and frequency of the first transition be $\nu_{Res1}$ where
Figure 3.3: Setup for saturated absorption spectroscopy.

Figure 3.4: Saturated absorption signal with 3 Lamb and 3 crossover dips for the $F = 4$ to $F' = 3, 4$ and 5 manifold of transitions.
\( \nu_{\text{Res1}} - \nu_1 \gg 5 \text{ MHz} \). Atoms in velocity classes centered in some small range (determined by the laser linewidth) around \(+ v_1\) and \(- v_1\) are then Doppler shifted into resonance with the left and right going beams, respectively. Now \( v_1 = \frac{1}{\varepsilon} (\nu_{\text{Res1}}/\nu_1 - 1) \). The detuning from transition is much more than the laser linewidth, and so atoms with velocities centered about \( + v_1 \) (\(- v_1\)) will absorb only from beams propagating in the \(-x(\pm x)\) direction. Since the left and right going beams are absorbed by atoms in completely different velocity classes, the spatial overlap of the beams does not affect the overall attenuation of the two.

At large detunings atoms with large velocities are Doppler shifted into resonance. The velocities along \( x \) follow a Gaussian distribution (figure 3.5). There are fewer atoms with large velocities, so the absorption will be smaller at larger detunings. As frequency is increased to \( \nu_2 \), closer to the first transition, atoms with lower velocities, \( \pm v_2 \), become resonant. Since there are more of these, the overall attenuation increases. Finally, as the laser frequency approaches within a linewidth of the 1st transition at \( \nu_{\text{Res1}} \), there begins to be overlap between the velocity classes of atoms which are resonant with the left and right going beams. This overlap increases until exactly on resonance the counter-propagating beams are absorbed by atoms in precisely the same velocity classes centered in some small range about \( v = 0 \). If the beam intensity is high enough to saturate the transition, then the overall attenuation of both beams decreases as the resonance frequency is approached since the effective total number of atoms absorbing from both beams is reduced. This produces a dip in the absorption profile of the photodetector, known as a Lamb dip (figure 3.4). Lamb dips occur at each hyperfine transition in the manifold. Similar considerations of absorption and velocity classes explain the presence of crossover dips which occur exactly midway in frequency between each pair of hyperfine transitions. Thus, scanning the laser frequency over the 3 hyperfine transitions in the manifolds from \( F = 4 \) and \( F = 3 \) produces a Doppler broadened absorption profile with six distinct dips (figure 3.4).

To frequency stabilize a laser requires a signal which varies with laser frequency, such as the side of the Doppler profile. As the laser frequency changes, the photode-
detector voltage increases or decreases depending upon the slope. The frequency of an external cavity laser can be tuned by changing the cavity length with a PZT. A positive voltage at the PZT can be made to increase the frequency while a negative voltage decreases it. An electronic feedback circuit which has the photodetector voltage as an input and which outputs a voltage to the PZT is able to stabilize the frequency. Figure 3.6 represents a portion of the side of the Doppler profile with one Lamb dip. The feedback circuit works in the following way. First, it compares the voltage from the photodetector, \( V_{pd} \), with a reference voltage, \( V_{Ref} \). The difference of these is called the error signal, \( V_{err} \), given by \( V_{err} = V_{pd} - V_{Ref} \). If \( V_{err} < 0 \), the circuit outputs a negative voltage to the PZT which reduces the laser frequency until \( V_{pd} = V_{Ref} \). If \( V_{err} > 0 \), a positive voltage is output to the PZT until, again,
When the laser is not stabilized, mechanical vibrations will cause it to fluctuate in frequency by several $\Gamma$ on 100 $\mu$sec time scales. Stabilization works as long as the error signal, $V_{err}$, changes sign when the laser frequency is scanned across the stabilization point within the range of normal frequency excursions due to vibration. Note that the center of the Lamb dip is an unsuitable point for stabilization in this scheme since if $V_{Ref}$ was set to the center of the Lamb dip, the sign of $V_{err}$ would not change as the frequency is scanned across the center. The overall stabilization method is shown in figure 3.7. Part of the light from the laser is sent through a saturated absorption setup and impinges upon a photodetector. The voltage from the detector goes into the feedback circuit which controls the laser frequency via the
Figure 3.7: Experimental setup for frequency stabilization to the side of the Doppler broadened profile

PZT. The repumper laser is controlled in this fashion.

The frequency of the cooling laser must be controlled with high accuracy and so it must be stabilized to one of the features in the absorption profile, the Lamb or crossover dips. But in the previous feedback scheme these were found to be unstable points. To circumvent this problem a lock-in amplifier is used to convert these features into an appropriate form for stabilization. First, light is double-passed through an AOM called the offset AOM and blue shifted by 160 MHz (figure 3.8). This frequency shifted beam is referred to as the cooling laser offset beam. This beam is split into two which are both double-passed through a Cesium vapor vessel just as in figure 3.8 and impinge upon a subtracting photodetection circuit.

When the laser frequency is scanned across hyperfine transitions, both paths produce a Doppler broadened absorption profile. The path which is retroreflected exhibits Lamb and crossover dips as in figure 3.4. The path which is not retroreflected has no such saturation features. Both profiles are subtracted at the photode-
Figure 3.8: Experimental setup for frequency stabilization to features in the saturated absorption profile

tection circuit and the Doppler broadened part vanishes leaving only the dips on a flat background. This is shown in figure 3.9 for transitions from the \( F = 4 \) hyperfine ground state.

Now the lock-in amplifier is employed. It puts out a time-varying voltage to the AOM which sinusoidally modulates the laser frequency at \( \omega_0 \) with small amplitude, \( \nu_{\text{mod}} \). For small amplitude oscillations in laser frequency the output signal from the subtraction circuit can be expanded in a Taylor series about some frequency, \( \alpha \), as

\[
f(\nu) = f(\alpha) + f'(\alpha)(\nu - \alpha) + \frac{1}{2} f''(\alpha)(\nu - \alpha)^2 + ...
\]  

(3.1)
Figure 3.9: The saturated absorption profile after subtraction of the Doppler broadened component (lower curve) and the output of the lock-in amplifier (upper curve).

The frequency modulation produced by the lock-in amplifier and AOM, then, has the form,

\[ \nu(t) = \alpha + \nu_{\text{mod}} \sin(\omega_0 t) \]  

(3.2)

Plugging this into the equation above and keeping terms to 1st order we have,

\[ f(\nu) \approx f(\alpha) + f'(\alpha) \nu_{\text{mod}} \sin(\omega_0 t). \]  

(3.3)

The lock-in amplifier then demodulates the signal using a sine wave at the modulation frequency, \( \omega_0 \), and measures the complex amplitude of the signal at this modulation frequency. The real part of this measurement is proportional to the slope, \( f'(\alpha) \).
If the laser frequency is scanned across the dips in the lower plot of figure 3.9 at a frequency much less than $\omega_0$, then the output signal from the lock-in amplifier will have the form of the upper plot in this figure. Note this function is sharply varying about the frequencies corresponding to the center of the dips, and represents a suitable signal for laser stabilization at these center frequencies in the feedback scheme already discussed. The output from the lock-in amplifier is put into an electronic feedback circuit which controls the PZT and thus, the laser frequency can now be stabilized to the center of a Lamb or crossover dip.

The cooling laser offset beam is then stabilized to the crossover peak between the $F = 4$ to $F' = 3$ and $F' = 4$ transitions, so the cooling laser itself is 160 MHz blue of this crossover peak (figure 3.10).

3.2.3 Injection-Locking

High power is needed in the MOT beams to load a large number of atoms. Thus, a powerful cooling laser is required. One of the disadvantages of the external cavity design is that the output power is much less than the rated power of the laser diode itself. This is a result of losses at the grating and a lowering of the maximum output power of the diode due to feedback. About seventy percent of the total laser power is reflected into the zero order at the grating. The other thirty percent is lost or diffracted into other orders. Fifteen percent of the power is diffracted back into the laser in the $-1$ order and together with the diode output makes up the intra-cavity power. Since the intra-cavity power cannot exceed the rating for the diode laser, it is not possible to operate the diode at full output power. However, a laser diode can be injection-locked by a small amount of light from a frequency stabilized "master" laser if their beams are properly mode matched. The injected laser (called the slave) matches the line-width and frequency of the master but is able to output virtually all of its rated power since the diffraction grating is eliminated and weak feedback is used. In our experiment the cooling laser is a 100 mW diode laser in an external cavity configuration which can output around 50 mW. Since the lasers are well mode matched and overlapped, two optical isolators which each provide 30 to 40 dB of
Figure 3.10: Frequency stabilization chart for lasers near the cooling transition. Lamb(crossover) dips are shown as solid(dashed) lines.
isolation are needed between them to protect the master from the slave output. The beam from the master laser is sent through the first isolator, mode matching optics and then into the side of the PBS cube at the output port of a second isolator (figure 3.11). At this location, the polarizations of the beams from the master and slave lasers are orthogonal. The isolator only isolates against light which has the same polarization as the beam from the slave laser. The beam from the master laser is reflected up through the isolator by the output PBS cube, and is then injected into the 200 mW slave laser. The slave laser is isolated from back-reflections off optical components after the second isolator. Since they have orthogonal polarizations, the output PBS cube of the second isolator also provides isolation of the master from the slave laser. At the right temperature and diode current the slave laser locks to the master and produces nearly 200 mW of output power at the same frequency as the master (figure 3.10). This high power laser is used to make the MOT and near resonance lattice beams.

3.3 MOT Optical Setup

3.3.1 AOM Double-Pass and Optical Fiber

Before breaking the slave cooling laser beam into the six MOT beams, it is beneficial to pass it through an optical fiber. The single-mode fiber transmits only the $TEM_{00}$ component of the beam and filters out higher order spatial modes. So, the output beam has a smoothly varying Gaussian profile which increases the density and uniformity of the MOT. The fiber also decouples alignment of all optics before and after itself which is a considerable time saver. To couple into the fiber, the laser is tightly focused onto the fiber facet with a ball lens. Back reflections from the facet which can cause severe feedback in the laser are alleviated by using index matching gel between the lens and facet. No such measures are necessary at the output facet. As mentioned earlier, to trap and cool atoms in 3D, it is necessary to change the frequency and intensity of the MOT beams dynamically. This must be done without changing the pointing of the beam from the slave cooling laser into the fiber and
is accomplished with an acousto-optic modulator in a double-pass geometry (figure 3.12). The collimated vertically polarized beam from the cooling laser is reflected off a polarization beam splitter (PBS) cube through an AOM. The first order which is blue shifted is bent parallel to the zero order by a lens placed its focal length from the AOM. This beam passes through a quarter wave plate, is focused onto a mirror where it is retroreflected and retraces its path back through the lens which collimates it. The beam then passes through the AOM again, is further blue shifted and deflected to overlap with the incoming beam. Since it is now horizontally polarized, the beam passes straight through the PBS cube. Immediately after the PBS cube
the beam is coupled into the optical fiber. Use of the lens to create optical path reversibility makes the beam pointing insensitive to frequency shifts of the AOM that change the deflection angle of the first order beam. The AOM can then be used to set the frequency of the MOT beams to the desired frequency around the $F = 4$ to $F' = 5$ transition.

### 3.3.2 MOT Components

After the fiber, the beam is spit into three linearly polarized beams of equal intensity which are individually passed through quarter-wave plates to produce circularly polarized light. They are then sent once through the chamber and afterwards pass through quarter-wave plates again before being retroreflected back through the chamber with the same circularity. This produces three pairs of counterpropagating beams of $\sigma+$ and $\sigma-$ light. A slight focusing is introduced to eliminate radiation pressure imbalance between counterpropagating beams. The MOT coils are centered about one pair of MOT beams and have a spacing of 12 cm between them. The coils
which have a diameter of 14 cm consist of 100 turns of insulated wire wrapped in a
water cooled copper mount. They can draw 6 amps indefinitely to produce magnetic
field gradients of 10 Gauss/cm during the loading phase.

3.4 Optical Lattices

3.4.1 Near Resonance Lattice

The near resonance lattice is formed by light taken from the cooling laser. The zeroth
order beam from the AOM double pass (figure 3.12), which is detuned 180 MHz red
of the $F = 4$ to $F' = 5$ transition, is deflected into a separate AOM which blue
shifts its frequency back toward resonance with the $F = 4$ to $F' = 5$ transition and
allows for fast control of the lattice beam intensity. Typical detunings of the near
resonance lattice range from 5 to 40 GHz. The beam is then passed through a single
mode fiber, and its output is collimated and split up to produce the lattice beams.
Since near resonance lattice intensities are low, of order $I_0$, only a few milliwatts of
light need be coupled through the fiber.

3.4.2 Far off-Resonance Lattice

In the far off-resonance lattice, which is typically detuned thousands of linewidths
off-resonance, high power is necessary to maintain potential depths of hundreds
of recoil energies, consistent with those of the near resonance lattice. Recall from
equation 2.7, the potential depth is proportional to the intensity divided by detuning.
The lattice light can be detuned arbitrarily as long as enough power is available.
Since the far off-resonance laser is detuned thousands of linewidths from resonance
as compared to 5 to 40 linewidths in the near resonance lattice, we require a high
power tunable laser system to make the far off-resonance lattice beams. The system
is designed around the SDL 8630 external cavity tunable diode laser (figure 3.13).
The device uses a tapered amplifier semiconductor chip as an active medium. Light
from the 5 $\mu$m rear facet is used with a diffraction grating in Littrow configuration to
provide feedback and wavelength selectivity. The output (front) facet which tapers
up to 150 $\mu$m can sustain powers of up to half a watt while maintaining single mode
operation. The output from the laser is near Gaussian and diffraction limited. Due to inadequate mechanical design, the laser has a poor linewidth which necessitates removal of the external cavity and use of another master laser with acceptable linewidth as a source for amplification. We drill a hole in the back of the laser enclosure to admit the master beam which is coupled into the 5 μm rear facet of the gain element. The input and output beams must be well mode matched for efficient coupling, but once this is achieved, the gain element acts as a single-pass amplifier of the master laser. The master laser is a second generation external cavity laser (appendix A) whose more monolithic construction offers better long term stability than earlier designs. This saves time since the laser need be re-aligned infrequently. Significant optical isolation is required between the master laser and amplifier chip. We use two OFR optical isolators (model IO-2-WIR2-HP) which each provide nearly 40 dB of isolation (figure 3.13). Inexpensive plano convex lenses are used to pass the master beam through the isolators as well as mode match it with the beam from the rear facet of the amplifier. A set of mirror mounts allow the beams to be spatially overlapped. High stability mirror mounts such as the Newport Ultima series (model U200-P) offer precision tuning and are necessary to maintain pointing stability over long times. Once it is set up and aligned, the system will operate reliably for weeks or months with very little maintenance.

Tunability is a critical issue for the master laser. Our only frequency standard is the saturated absorption signal. This means the master laser must start on transition and be tuned smoothly, without mode hopping, to its final frequency. The amount of tuning is measured by a laser spectrum analyzer with a 2 GHz free spectral range. The far off-resonance lattice is, generally, detuned 10 to 40 GHz from transition. Our external cavity lasers can be tuned, by varying cavity length alone, only up to 3 GHz before mode hopping. However, with an appropriate combination of cavity length and diode current tuning, the laser frequency can be smoothly scanned more than 40 GHz. The output beam of the high power laser system is sent through an optical isolator to protect it from unwanted feedback and then through and AOM to provide intensity control of the far off-resonance lattice beams.
3.5 Diagnostic Tools

3.5.1 Time-of-Flight Analysis

The primary diagnostic tool in these experiments is Time-of-Flight (TOF) Analysis [3] which allows one to extrapolate the temperature of atoms at their time of release from an optical lattice. Temperature is measured along one direction only, the gravitational axis. In this technique the atoms are released from the lattice and fall. The initial size of the sample is small, but since it has a nonzero temperature, it will expand in all directions. After about 98 ms, the atoms pass through a thin probe beam. A photodetection circuit measures absorption of the probe beam by the expanded atomic cloud as it falls through the probe and produces a distribution of atomic arrival times. This distribution can be converted into a distribution of momentum for the atoms in the lattice at the time of their release. The temperature
of the atoms at release is extracted from the momentum distribution.

Starting with a normalized Gaussian momentum probability distribution,

\[ \Pi(p) = \frac{1}{\sqrt{2\pi \cdot M k_B T}} e^{-\frac{p^2}{2Mk_BT}} \]  

(3.4)

the temperature can be expressed in terms of the variance, \((\sigma_p)^2\), of the distribution by,

\[ \sigma_p^2 = M k_B T \]  

(3.5)

Here \(T\) is the temperature and \(M\) is atomic mass and \(p\) is the momentum of an atom upon release from the lattice along the gravitational axis. This can be converted into a probability distribution of arrival times, \(\Pi(t)\), by using [26],

\[ \Pi(t) = \Pi(p) \left| \frac{dp}{dt} \right| \]  

(3.6)

The relationship between the initial momentum, \(p\), of an atom and its arrival time, \(t\), at the probe beam is given in this case by the equation for a freely falling particle,

\[ p = Mg \left( \frac{t^2 - t_0^2}{2t} \right) \]  

(3.7)

\[ = Mg\Delta t \left( 1 - \frac{1}{2} \frac{\Delta t}{t} \right) \]  

(3.8)

where \(t_0 = \sqrt{2d/g}\) is the arrival time of atoms initially at rest, \(d\) is the distance fallen, \(g\) is the acceleration due to gravity, and \(\Delta t = t - t_0\). Taking the derivative of \(p\) with respect to \(t\) and plugging into equation 3.6, we find the probability distribution for atomic arrival times,

\[ \Pi(t) = \frac{1}{\sqrt{\frac{2\pi k_B T}{Mg^2}}} \left( 1 - \frac{\Delta t}{t} + \frac{1}{2} \left( \frac{\Delta t}{t} \right)^2 \right) e^{-\frac{(\Delta t)^2}{2(k_B T/Mg^2)}} \left( 1 - \frac{1}{2} \frac{\Delta t}{t} \right)^2 \]  

(3.9)

If the spread in arrival times is small compared to the fall time, then \(\Delta t/t\) is small and can be neglected. For this case the population distribution in time has the form of a Gaussian and the temperature of the atoms at the time of release can be written in terms of its variance, \((\sigma_t)^2\),
The total integrated area of the Gaussian TOF signal is proportional to the number of atoms in the sample.

Experimentally, the probability distribution of arrival times is found from the absorption of light by atoms falling through a probe beam. The probe beam is sent through two pairs of anamorphic prisms which expand it along one axis, then collimated, and then a thin slit selects the central portion of the elongated beam. The beam is split into two, a signal and a reference beam. The signal beam is focused by a cylindrical lens along its short axis to a 20 µm thick waist directly beneath the atoms; it is about one inch wide along the other axis. Atoms fall through the signal beam only, and then, signal and reference beams are gathered onto a photodiode subtraction circuit which measures differential absorption between them and rejects their common mode power fluctuations. Not surprisingly, signal to noise is a problem since the maximum power absorbed by a few million atoms in the signal beam is of order 10 nW. It is generally necessary to average the signal over many drops. Figure 3.14 shows TOF distributions with overlayed Gaussian fits for atoms released from 3D optical molasses both for a single release and an average over twenty releases.

Since it is important to know temperature as accurately as possible, corrections to the temperature found using TOF analysis should be considered. The most important is due to finite MOT size. An initial (assumed) Gaussian spatial distribution of the MOT gives rise to a Gaussian distribution of arrival times even at zero temperature. This function is convolved with the distribution of arrival times due to nonzero temperature at the probe beam and results in an artificially high apparent kinetic temperature. From the convolution theorem applied to Gaussians it is easy to see that the variance of the convolved function equals the sum of the variances of the original functions. Since temperature is proportional to the variance, the contribution due to finite MOT size is additive and can be accounted for by subtracting a constant temperature. For a MOT with a FWHM of 400 µm in its spatial den-
Figure 3.14: TOF distributions for atoms released from 3D optical molasses for a single release (upper curve) and for an average of 20 releases (lower curve). Gaussian fits to the data are indicated by solid lines.
sity distribution, the correction to temperature is around 50 nK. Another correction due to finite probe thickness, can be estimated similarly by assuming a Gaussian intensity profile for the probe. The 20 \( \mu \text{m} \) thick probe causes a negligible correction on the order of 1 nK. Finally, corrections due to the time domain response of the photodetector electronics with its 30 \( \mu \text{s} \) time constant are also negligible.

### 3.5.2 Zeeman Analysis

In addition to temperature it is also useful to have information about the internal state of the atoms when they are released from the lattice. Zeeman state analysis involves finding the population distribution of the sample of atoms over the nine Zeeman sublevels in the \( F = 4 \) hyperfine ground state, and is carried out much like a Stern-Gerlach experiment. A strong magnetic field gradient is used to separate out the arrival times of atoms in the nine magnetic sub-levels at the probe beam. The force on an atom with a fixed magnetic moment, \( \hat{\mu} \), in a magnetic field is,

\[
\mathbf{F} = \nabla (\hat{\mu} \cdot \mathbf{B}),
\]

where \( \mathbf{B} \) is the magnetic induction. When the atoms are released from the lattice, a 1 Gauss uniform magnetic “alignment” field is applied along the lattice quantization axis to prevent precession of their magnetic moments due to accidental stray fields, and to preserve their projections of angular momentum, \( m\hbar \), onto this axis. The atoms fall for 24 ms and then the MOT coils are switched on for 50 ms. The MOT coils produce a magnetic field larger than 1 Gauss, which is not along the direction of the alignment field and so the total magnetic field begins to change direction. If the rate of this change is slow, compared to the Larmor precession frequency for the atomic magnetic dipole moment, then the dipole moment will adiabatically follow the changing magnetic field. In that case the projections, \( m\hbar \), onto \( \mathbf{B} \) are preserved, and the force depends upon the gradient of only the magnitude of the magnetic field, and is given by,

\[
\mathbf{F} = -(g_F \mu_B m) \nabla |\mathbf{B}|.
\]
Twenty-four milliseconds after release from the lattice, the atoms have fallen into a region where $\nabla|B|$ is primarily along the gravitational axis, and has a value of about 11 Gauss/cm. The nine Zeeman sublevels with different projections of angular momentum, experience different forces and are separated spatially along the $z$-axis during the 50 ms period for which the MOT coils are on. This separates their arrival times at the probe beam. The TOF distribution can be fit by nine Gaussians, but the different forces experienced by atoms that follow different trajectories skew the apparent temperature and integrated area of the individual peaks. By carefully modeling atomic trajectories for our particular experiment, we find correction factors for these fits and are able to extract reliable values for the population and temperature of each sublevel.
Figure 3.15: TOF distribution with nine peaks for atoms released from a near resonance lattice with a magnetic field gradient present.
CHAPTER 4

EXPERIMENTS IN A ONE-DIMENSIONAL LATTICE

4.1 The 1D Lin⊥Lin Optical Lattice

4.1.1 Introduction

This chapter describes work carried out in a 1D Lin⊥Lin lattice which had as one of its goals assessment of the feasibility of quantum state preparation and control in such a lattice. In particular, it was conjectured that resolved-sideband Raman cooling might be possible in this system [11] and, though this turned out not to be the case, much of what was learned was applicable to the sideband cooling experiment later carried out in a two-dimensional lattice, which was ultimately successful.

The sideband cooling experiment involves cooling all of the atoms in the lattice into the lowest vibrational level in a particular Zeeman sublevel into which they have all been optically pumped. Therefore, it is advantageous to begin this experiment with as many atoms as possible in the appropriate magnetic sublevel and with a low initial mean vibrational excitation, $\bar{v}$. It is possible to load directly from 3D molasses into a far off-resonance lattice, but the absence of rapid laser cooling in the FRL means that the mean vibrational excitation would be high. In principle, the lattice potential depth could be lowered to allow atoms in excited states to escape while retaining only atoms in low lying vibrational levels [27, 28], but this would come at the cost of losing most of the atoms. In practice, since the signal produced by the atoms is weak, it is undesirable to work with only a small subset of those cooled in 3D molasses. In a near resonance lattice the same high photon scattering rate which makes it unamenable to sideband cooling, allows for laser cooling which causes the atoms to rapidly become trapped and localized around the bottom of individual potential wells with low mean vibrational excitation. We
were motivated, then, to attempt transfer of atoms from a near resonance into a far off-resonance lattice while maintaining the strong localization produced in the NRL.

For transfer, in terms of the goal of sideband cooling, the primary considerations were whether the low $\overline{n}$ from the NRL could be preserved and whether the magnetic sublevel distribution could be maintained. Also important was to determine the rate of vibrational heating in the FRL. For sideband cooling to work, this rate must be much less than the sideband cooling rate. It should be noted that for experiments in 1D lattices, producing a sample with very low $\overline{n}$ in the FRL may be even more important, since for these experiments sideband cooling with our scheme is not possible, and $\overline{n}$ must be reduced by selecting only the subset of atoms which are already in the lowest lying vibrational level.

Finally, the reasons for working in a 1D lin $\perp$ lin optical lattice are many. From a theoretical standpoint it is a well studied system for which band structure calculations are possible. In the case where atoms populate vibrational levels near the top of the potential well, band theory offers a better description of the lattice than an anharmonic oscillator. Experimentally, a 1D lattice requires fewer optical elements and is easier to set up and align than higher dimensional lattices.

4.1.2 Optical Components of the Lattice

The superimposed near resonance and far off-resonance 1D lin $\perp$ lin optical lattices are made in the following way. Near resonance and far off-resonance beams travel through separate AOM's which provide independent intensity control of each, then through half-wave plates and into separate faces of a polarization beam splitter cube (figure 4.1). The half-wave plates fix the splitting ratio of the cube at 50% for both beams while the PBS cube produces, for near resonance and far off-resonance light, beams of orthogonal polarization which are made to counter-propagate by mirrors after the cube. Copropagating near resonance and far off-resonance beams clearly have orthogonal polarizations. This is allowable due to the high degree of symmetry of this lattice and since this is so, the near resonance and far off-resonance light can
be combined without any loss which is not generally the case for an arbitrary optical lattice. The region of overlap of the counter-propagating beams after the PBS cube is the optical lattice. The section of the lattice which intersects the MOT in the vacuum chamber is aligned with the gravitational axis to better than 20 mrad. The length of an arm of the lattice is defined as the distance from the PBS cube to the MOT position along either of the two possible paths. As will be explained, it is important for good transfer that both arms have the same length. Placing lattice mirrors on translation stages allows for fine tuning of the length difference between the two arms.

The first method used for aligning the lattice is as follows. A lens is inserted into one arm of the lattice in such a way as to focus near resonance light at the MOT position without translating the beam transversely. A mirror after the lens allows the focused beam to be steered to the center of the MOT where it pushes atoms away and creates a hole which can be viewed with a CCD camera. This centers the near resonance lattice beam on the MOT very well. A plumb bob determines beam verticality. The lens is then removed and lattice mirrors in the other arm are used to overlap the counter-propagating beam on top of the former. Finally, mirrors in the far off-resonance light path external to the lattice are employed to overlap the FRL on the NRL. Reliable calculations of potential depth and scattering rate in the FRL are very important and rely upon having the FRL beams well centered on the MOT so that the intensity at the MOT position is well known.

4.1.3 Far off-Resonance Lattice Intensity Profile

Since knowledge of the intensity experienced by atoms in the far off-resonance lattice is vital, the intensity profile of the far off-resonance lattice beams is of great concern. For example in the extreme case, if the transverse extent of the lattice beams is too small, the atoms, which are not cooled transverse to the 1D lattice, can quickly leave the lattice volume before the experiment is performed. In most cases it is ideal for all atoms to experience the same lattice intensity over the course of the experiment. This suggests a “top hat” or uniform intensity profile is best. Such profiles can be
produced but often with some difficulty and at the price of reduced lattice beam intensity. Laser sources commonly output Gaussian intensity profiles. If uniform intensity is required in the lattice, the FWHM of the beam can be made much larger than the sample of atoms so that intensity variations will be small across the sample.

For our initial experiment involving adiabatic transfer into a FRL and measurement of $\bar{n}$, it was not necessary to have highly uniform beams since the effects of nonuniformity for that experiment could be handled by calculation. Much of the data was taken in a FRL with a Gaussian beam profile with FWHM of 3.3 mm. This leads to a half-width at ninety percent of maximum of 585 $\mu$m. The temperature of atoms in 3D molasses is around 3.5 $\mu$K. The atoms are further heated transversely while being held in the NRL. By the time the atoms are transferred to the FRL.
they can have RMS speeds of around 20 mm/sec. So over the course of 20 ms in
the FRL and considering a 400 μm MOT size, a significant fraction of atoms has
moved into a region of the lattice where the intensity has fallen by more than 10% as compared to the on axis intensity. The effect of slowly moving into a region of
lower intensity is that the atom experiences a potential well which is adiabatically
lowered in depth. Since the process is adiabatic, it does not change the distribution
over vibrational levels of such atoms and, thus, the distribution is independent of
the position of an atom in the lattice. This fact allows the calculation of \( \bar{n} \) for a
sample of atoms in such an inhomogeneous lattice as long as the Gaussian intensity
profile is known as well as the spatial distribution of atoms in the lattice.

4.2 Adiabatic Transfer into a Far off-Resonance Lattice

4.2.1 Adiabatic Release and Recapture

To transfer atoms from a near resonance to a far off-resonance lattice, it is necessary
to have precise control of the intensity of both sets of lattice beams and to be able
to independently vary the depth of the lattices in time. The intensity of a lattice is
controlled by an AOM which deflects more power into the lattice beams as it receives
more power from an RF source. The deflection efficiency of an AOM is nonlinear in
the applied RF power. Once this efficiency is characterized, an arbitrary function
generator can be used to control the amplitude of the RF signal at the AOM to
produce any desired lattice intensity waveforms.

During the development of the transfer technique many transfer intensity wave­
forms and timing sequences were tried. Only two will be discussed here. The first
uses adiabatic release from the near resonance lattice followed by adiabatic capture in the far off-resonance lattice. The intensity of the near resonance lattice is
adiabatically lowered according to the equation,

\[
I_f(t) = \frac{I_i(t = 0)}{(1 + \xi t)^2}
\]  

(4.1)

where \( I(t) \) is the intensity at time \( t \), \( I' \) is the initial intensity and \( \xi \) is the adiabatic
release parameter in units of inverse time. As the lattice depth is lowered, the tightly
bound atoms slowly become free particles with momentum close to zero. At this point the far off-resonance lattice is adiabatically turned on with a waveform which is the time reversal of the one above. In principle, the free particle states should become tightly bound states in the far off-resonance lattice [29]. An advantage of this type of transfer is that it should work well even if the NRL and FRL are not topographically similar or, if they are similar, even if their potential wells are not spatially overlapped. In practice in our setup this technique does not work well. During the time when the atoms are in free states, gravity accelerates them so that the FRL is left to adiabatically capture free states of non-zero momentum which provides an increase in the mean vibrational excitation over that of the NRL. Further, while the atoms are in free states, there is very little lattice light and the quantization axis is not well defined. Then the magnetic moment of an atom can begin to precess about random background magnetic fields which means that the Zeeman sublevel distribution of the NRL will not be preserved upon recapture in the FRL. Both of these problems can be dealt with. Release and recapture can be carried out in a freely falling frame to eliminate gravitational effects and a magnetic field can be applied along the lattice quantization axis during transfer to preserve the distribution over magnetic sublevels. But for reasons of experimental ease a different transfer method is used.

4.2.2 Tight-Binding Transfer

The second transfer method involves maintaining tight binding of the atoms throughout the transfer. Since the intensity of the lattice beams can be well calibrated and controlled, it is possible to sustain the apparent potential depth experienced by the atom during transfer. This is accomplished by linearly changing the intensities of the NRL down and FRL up in such a way that the sum of the depths of the two potentials is approximately constant throughout the transfer. The parameters of the transfer including slope, duration and relative starting time of the linear intensity change for both lattices were adjusted extensively to optimize the transfer with regards the efficiency (percentage of atoms transferred) and minimization of tem-
perature. Eventually, it was realized that the transfer temperature is very sensitive to the spatial overlap of the topographically similar potentials of the near and far off-resonance lattices. This is not surprising. If the potential wells of the NRL and FRL are not overlapped, then atoms which begin in regions of low potential around the bottom of the NRL wells find themselves, upon transfer, in regions of high potential along the sides of the FRL wells. This increased potential energy translates into an increased mean vibrational excitation upon equilibration in the FRL. So a method was devised to insure spatial overlap of these potentials.

4.2.3 Overlapping NRL and FRL Potentials

The detuning between the NRL and FRL is of order 10 GHz. The wavelength change this represents is tiny compared to the wavelength, \( \lambda \), of the laser light. There is, therefore, very little difference in periodicity between the near and far off-resonance lattices. The lattice is a standing wave set up by two counterpropagating beams. Consider a complete circuit around the lattice starting at the PBS cube and returning to it (figure 4.1). At the midpoint of this path, called the midpoint of the lattice, the standing wave, independent of its wavelength, is constrained to have a node. So at this location the NRL and FRL potential wells are overlapped for any FRL detuning. When the two arms of the lattice have equal length, the MOT sits at the midpoint of the lattice and the transfer temperature into the FRL is minimized for any detuning of the FRL, since changing the FRL frequency does not significantly alter the overlap of the potential wells across the MOT. If the path lengths of the two arms of the lattice differ, this is equivalent to moving the MOT away from the midpoint of the lattice. Far enough from the midpoint, the difference in periodicity of the two lattices causes their potential wells to become significantly mis-aligned, and this shows up in the transfer temperature. For a detuning of about \( 3831 \Gamma \) (20 GHz) in the far off-resonance lattice, the spatial beat period between near and far off-resonance potential wells is around 1.5 cm. This suggests a method for precision zeroing any length difference between the two arms of the lattice. Varying the FRL frequency changes the spatial overlap of the potential wells if the MOT is not at the
midpoint of the lattice. A plot of transfer temperature versus FRL detuning reveals a function which is sharply peaked in temperature when the periodic potentials of the NRL and FRL are completely out of phase (figure 4.2a). The spacing between these temperature maxima indicates how much frequency change of the FRL is needed to shift its phase by $2\pi$ with respect to the NRL and from this, the distance of the MOT from the midpoint or length difference of the lattice arms can be determined and is given by,

$$d = \frac{\lambda^2 \nu^2}{c|\Delta \nu|}$$

where $\Delta \nu$ is the frequency shift between adjacent temperature maxima. For the data in figure 4.2a, $\Delta \nu = 3.1$ GHz which means that the path lengths of the two lattice arms differ by about 9.7 cm. This is corrected by translating a mirror in one arm of the lattice in the appropriate direction. Then the potential wells are overlapped and there is no frequency dependence to the transfer temperature (figure 4.2b).

4.2.4 Results of Transfer

After the problem of potential well overlap was solved, this tight binding transfer technique was used exclusively. Figure 4.3 shows TOF distributions for atoms released from near resonance and far off-resonance lattices, which are separated in their arrival times at the probe beam by the time for which the atoms are held in the FRL, 1.3 ms and 20 ms for the data in figure 4.3. Fitting Gaussians to these peaks reveals a 90% transfer efficiency. The signal is still quite visible at 20 ms, but atoms are clearly being lost. This loss rate from the FRL is well fit by a decaying exponential. The major source of loss is due to atoms being heated by photon scattering and escaping from their potential wells. Other sources of loss include decay to the $F = 3$ hyperfine ground state in which case an atom is no longer bound in the lattice, and at long times escape from the lattice volume by transverse motion. The distribution over magnetic sublevels is also well preserved during transfer using this technique. Figure 4.4 shows this qualitatively for two different preparations in
Figure 4.2: Deviation from expected temperature versus detuning in a far off-resonance lattice (a) before and (b) after the path difference between the two arms of the 1D lattice has been eliminated.
the near resonance lattice.

4.3 State Preparation in a Far off-Resonance Lattice

4.3.1 Experimental Procedure

One of our early experiments involved the study of transfer into and properties of a 1D lin \perp lin far off-resonance lattice and had as one goal assessment of the feasibility of resolved-sideband Raman cooling in far off-resonance lattices. Results from this experiment were published in [12]; it is carried out as follows. Atoms are loaded into a 1D lin \perp lin NRL from a magneto-optic trap and 3D molasses for about 1 ms where they are cooled and localized around the well bottoms. They are then transferred into a 1D lin \perp lin FRL as described before by a linear turn-off of the NRL and turn-on of the FRL which takes place over 200 \mu s. The atoms are held in the FRL for a time and released by instantaneously extinguishing the light.

4.3.2 Preservation of $n$ during Transfer

One of the important quantities to calculate is the mean vibrational excitation, $\bar{n}$. For atoms released from the lattice, this is extracted from their RMS momentum, $\Delta P$, which is found experimentally from a Gaussian fit to the TOF distribution of the sample (section 3.5.1). The mean vibrational excitation, $\bar{n}$, is calculated from $\Delta P$ as described in section 2.3.2, but the procedure is complicated by the Zeeman sublevel structure of the atoms as well as lattice beam inhomogeneity. Assuming that $\bar{n}$ is the same for each magnetic sublevel and independent of local lattice intensity, its value is calculated as described in [12].

Near resonance lattice parameters ($I = 2.5I_0$ and $\Delta = -23\Gamma$) were selected to prepare the atoms with low $\bar{n}$ and were maintained throughout the experiment. Far off-resonance lattice parameters were varied. In each case the timing of the transfer between lattices was optimized to minimize $\bar{n}$ in the far off-resonance lattice. The mean excitation inferred from the TOF distribution of the near resonance lattice is $\bar{n} = 0.37 \pm 0.03$. This corresponds to a total ground state population, $\Pi_{n=0} = 0.73$, 


Figure 4.3: TOF distributions for transfer from a near resonance (dashed line) to a far off-resonance (solid line) lattice for atoms which are held in the far off-resonance lattice for (a) 1.3 ms and (b) 20 ms.
Figure 4.4: TOF distributions under Zeeman state analysis in a NRL (upper curves in (a) and (b) ) and after transfer to a FRL (lower curves in (a) and (b) ).
which means that 73% of atoms are in the ground vibrational level. Note that 
\( \Pi_n \) is the total fraction of all atoms, irrespective of Zeeman sublevel, which are 
found in the ground state of their respective potential wells. Magnetic sublevel 
analysis (figure 4.5) reveals about 75% of atoms are in the stretched \( |4, m = \pm 4) \) 
states. The population imbalance between \( |4, m = +4) \) and \( |4, m = -4) \) is likely 
caused by a nonzero magnetic field along the lattice quantization axis. The lowest 
bound states in the lattice are the ground vibrational levels of the stretched states 
\( |4, m = \pm 4, n = 0) \). From the population in \( |4, m = \pm 4) \) and \( \Pi_{n=0} \), the population 
in the lowest bound states in the lattice is found to be about 55%. This is consistent 
with a more careful determination of population in the lowest bound states made 
later [15].

Figure 4.6 shows the mean excitation \( \bar{n} \) of atoms transferred into the FRL for
Figure 4.6: Mean vibrational excitation versus peak-peak modulation depth, $U_{pp}^A$, of the diabatic potential associated with state $|4, m = 4\rangle$ in the far off-resonance lattice. Solid (open) symbols represent data taken for lattices detuned $3831\Gamma(1916\Gamma)$ red of the cooling transition. Squares (circles) indicate data taken for atoms held in the far off-resonance lattice for 0.1 ms (20 ms).

a variety of potential depths, detunings and hold times. For very short hold times before the atoms can heat significantly in the FRL, we find $\bar{n} = 0.34 \pm 0.03$. This demonstrates that $\bar{n}$ can be preserved during transfer. Thus, the FRL can be efficiently loaded and achieve localization typical of a NRL. Magnetic sublevel analysis (figure 4.5) demonstrates that about 55% of atoms are in the states $|4, m = \pm 4\rangle$ in the FRL. So the magnetic sublevel distribution is well, though not perfectly, preserved after transfer.
4.3.3 Heating Rate in the FRL

The rate of increase of $\bar{n}$ in the far off-resonance lattice is referred to as the heating rate. Figure 4.7 plots $\bar{n}$ versus the time for which atoms are held in the far off-resonance lattice at various detunings and potential depths. There is clearly an increasing trend for $\bar{n}$ with time. One can calculate the rate of increase of $\bar{n}$ due to photon scattering for atoms in potential wells of the $|4, m = \pm 4\rangle$ states by solving rate equations for the vibrational populations. These calculations are plotted as lines on figure 4.7 and are found for scattering rates consistent with the plotted data which range from 500 to 2000 s$^{-1}$. At short times the data roughly follows this calculated linear increase but departs from it at long times. Surprisingly, the temperature at long times appears to depend upon the lattice depth and not the photon scattering rate as expected. Consider the two sets of data in figure 4.7 taken at the same lattice depth, $U_{pp} = 105E_R$, plotted as circles. As expected, at short times, less than 5 ms, in the far off-resonance lattice the set with lower detuning and higher scattering rate (open circles) has a greater heating rate. However, at long times, heating rates for both sets are indistinguishable. The data taken for atoms held in the far off-resonance lattice for long times, 20 ms, in figure 4.6 (circles) indicates a linear increase of $\bar{n}$ with lattice depth.

As the hold-time increases, so does uncertainty in $\bar{n}$. This occurs because atoms are heated into higher vibrational levels in their potential wells and also move into regions of shallower lattice depth. When this happens, the anharmonic-oscillator model from which $\bar{n}$ is calculated becomes a poor approximation. Therefore, reliable data cannot be taken for long enough hold-times to determine whether $\bar{n}$ reaches a steady-state value. Since the heating rate due to photon scattering should always be present, if $\bar{n}$ were to come to steady-state, that might indicate the presence of a cooling mechanism in the FRL [30].

4.3.4 Feasibility of Sideband Cooling

We are now in a position to assess the prospects for resolved-sideband Raman cooling in a far off-resonance lattice. Sideband cooling can remove a quantum of vibrational
energy from a bound atom on time scales of the order of a vibrational period. The heating rate is maximum immediately after transfer into the FRL. For the data in figure 4.8 at short times, it ranges from 20 to 100 s\(^{-1}\).

The atomic oscillation frequencies in potential wells of depths \(U_{pp} = 105E_R\) and \(U_{pp} = 209E_R\) are 42 kHz and 60 kHz, respectively. Taking the ratio of heating rate to oscillation frequency in the extreme cases gives \(0.48 \times 10^{-3}\) and \(1.7 \times 10^{-3}\). So the heating rate can be as little as a thousandth of the expected sideband cooling rate, and this indicates the prospects for sideband cooling in a FRL are good, at least insofar as consideration of rates is concerned.
4.4 Enhanced State Preparation in a Near Resonance Lattice

4.4.1 Introduction

Now that we have seen that transfer from a near resonance into a far off-resonance lattice can preserve mean vibrational excitation and distribution over magnetic sublevels, we are motivated to optimize the preparation of atoms in the near resonance lattice to improve experiments in the far off-resonance lattice. For example, since the goal of sideband cooling is to prepare all atoms in a single stretched state in the ground vibrational level, it is reasonable to prepare the maximum number of atoms in this stretched state in the near resonance lattice with $\bar{n}$ as low as possible. As mentioned in [15], for many experiments it is helpful to prepare atoms in the near resonance lattice in a state which is close to a pure state. Preparations which involve increasing population in one stretched state and lowering $\bar{n}$ can be enhanced by the application of weak magnetic fields. This section is a brief summary of experiments and findings involving the application of weak magnetic fields in 1D lin $\perp$ lin optical lattices. A more detailed discussion can be found in [15].

4.4.2 Weak Magnetic Fields in the NRL

As discussed in section 2.2, the effect of applying a longitudinal magnetic field along the lattice quantization axis is to Zeeman shift the diabatic potentials in energy with respect to each other according to,

$$\delta E_m = (g_F \mu_B m) B_z \quad (4.2)$$

where $\delta E_m$ and $m$ are the energy shift and quantum number, respectively, of the $m^{th}$ Zeeman sublevel and $B_z$ is the longitudinal magnetic field. When the adiabatic potentials are considered, the effect of a longitudinal field on the lowest potential is to shift the potential wells associated with the two stretched states in energy with respect to each other. This, effectively, makes one deeper with more bound states and the other shallower with fewer. The tendency is for population to build up in the deeper potential well as atoms are pumped into it from the shallower
wells [14]. Atoms bound in the deeper well are in eigenstates which reflect, almost entirely, the character of the stretched state associated with that well. Figure 4.8 shows the adiabatic potentials for an \( F = 4 \) atom in a 1D lin \( \perp \) lin optical lattice composed of counterpropagating beams aligned along the z-axis. For the field-free case (figure 4.8a) the lowest potential wells associated with the stretched states are symmetric. Application of a weak longitudinal magnetic field, \( B_z \), breaks this symmetry 4.8b. Transverse magnetic fields introduce off-diagonal elements into the light shift operator which couple magnetic sublevels that differ by \( m = \pm 1 \). These couplings modify the adiabatic potential as shown in 4.8c and cause mixing among magnetic sublevels to occur in regions of the lattice closer to the well bottom than in the case of no background field. A cooling mechanism described by Deutsch et al. [31] relies upon coherent mixing of magnetic sublevels and can be expected to produce lower temperatures when mixing is increased by the presence of transverse fields.

We explore the effects of weak magnetic fields on a near resonance lattice in [15]. The experiment is described in the following sections and its results are consistent with the behavior discussed above.

### 4.4.3 Experimental Procedure

Atoms are loaded from a magneto-optic trap and 3D molasses into a 1D lin \( \perp \) lin near resonance lattice made of counterpropagating beams aligned along the z-axis and detuned about 20 linewidths red of the \( F = 4 \) to \( F' = 5 \) hyperfine transition. Separate coils in the xy and yz-planes are used to produce magnetic fields along the z (longitudinal) and x (transverse) axes while atoms are held in the NRL. Then the near resonance light is instantaneously turned off and the temperature and magnetic sublevel distribution of the sample is measured.

### 4.4.4 Longitudinal Magnetic Fields

Results for the application of longitudinal magnetic fields are given in figure 4.9a. The plot shows the populations, \( \Pi_{\pm 4} \) in either state \( \ket{4, m = \pm 4} \) as a function of
Figure 4.8: Adiabatic potentials in a 1D linear optical lattice of detuning, $\Delta = -20\Gamma$, and single beam light shift, $U_1 = 45E_R$ for (a) no magnetic field, (b) a longitudinal magnetic field of 30 mG applied along the z-axis, and (c) a transverse magnetic field of 30 mG applied along the x-axis.
the longitudinal field. Longitudinal field is converted into the Zeeman shift of the \(|4, m = \pm 4\rangle\) states, \(\delta E_{\pm 4}\), and scaled by the peak-peak modulation depth of the corresponding diabatic potentials \(U_{pp}^4\) and \(U_{pp}^{-4}\) for data in the plot. The maximum population in either stretched state occurs when the Zeeman shift is about half the peak-peak potential depth. For higher fields the disparity between the depths of the potential wells in figure 4.8b becomes great and the anticrossings vanish. Then the population in the deeper potential well begins to fall off. The highest stretched state population found during the experiment, was 0.76. Lines on the plot indicate quantum Monte Carlo wave-function (QMCWF) simulations of population in the stretched states for the given lattice parameters [32]. The data is in good agreement with these simulations and any differences can be explained by imperfect control of magnetic fields at the 10 mG level.

From the momentum spread of atoms in the two stretched states, which we find using Zeeman state analysis, we extract the vibrational temperature and population distribution over vibrational levels. Figure 4.9b shows the ground state populations, \(\Pi_{\pm 40}\), in potential wells associated with the two stretched states, plotted against the longitudinal magnetic field which has been converted into Zeeman shift. The maximum ground state population is achieved for Zeeman shifts of about one quarter of the modulation depth of the associated diabatic potentials. This is not the same value which maximizes the total population in either extreme magnetic sublevel, \(\Pi_{\pm 4}\). The longitudinal magnetic field also causes heating which decreases the population in the ground vibrational level. The Zeeman shift which maximizes the ground state population in potential wells associated with either stretched state, \(\Pi_{\pm 40}\), represents a compromise between increasing \(\Pi_{\pm 4}\) while not increasing \(\Pi\) too much. Experimentally, the highest ground state population found is \(\Pi_{\pm 40} = 0.28 \pm 0.03\) which is a 33% increase over the zero-field value of 0.21 \(\pm 0.03\).

QMCWF simulations are also shown in figure 4.9b [32] and agree with the trend of experimental data but indicate higher peak populations than actually found. Disagreement is, again, probably due to background fields and imperfections in the lattice beams.
Figure 4.9: (a) Total population in magnetic sublevels $|4, m = 4\rangle$ (open boxes) and $|4, m = -4\rangle$ (filled boxes) versus applied longitudinal field in units of $\delta E_{\text{z}} / U_{pp}$ for lattice detuning $\Delta = -20\Gamma$ and modulation depth $U_{pp}^4 = 135E_R$. Lines represent QMCWF simulations corresponding to the data. (b) Vibrational ground state population of potential wells with $|4, m = 4\rangle$ (open boxes) and $|4, m = -4\rangle$ (filled boxes) character versus applied longitudinal field in units of $\delta E_{\text{z}} / U_{pp}^4$ for lattice parameters $\Delta = -20\Gamma$ and $U_{pp}^4 = 70E_R$. Lines represent QMCWF simulations.
4.4.5 Transverse Magnetic Fields

In figure 4.10 is plotted kinetic temperature versus transverse magnetic field. The lower set is for zero longitudinal field while the upper set is for $B_z = 90$ mG. Both sets reveal a reduction in kinetic temperature for the application of weak transverse fields as anticipated. QMCWF simulations for the two data sets show general agreement with the data [32]. Discrepancies here are easily attributed to uncontrolled background fields of a few 10's of milliGauss. We are primarily interested in the distribution over vibrational levels in the lowest adiabatic potential. But in this case, the strong mixing among magnetic sublevels in the various potential wells precludes Zeeman analysis as used before to find the momentum spread associated with atoms in a particular potential well. This is because in the presence of strong mixing, atoms in a particular potential well can no longer be associated with a single magnetic sublevel. This is not a problem, however, for QMCWF simulations which can be used to calculate $\Pi_{\pm 40}$. Such simulations suggest optimal magnetic field and lattice parameters. For $B_x \approx B_z \approx 25$ mG and $U_{pp} = 70E_R$, simulations show the ground state population increases by 80% over the value of data taken for $B_x \approx B_z \approx 0$, from 0.25 to 0.45.
Figure 4.10: Kinetic temperature versus transverse magnetic field applied along the x-axis, for static longitudinal magnetic fields $B_z = 0 \text{ mG}$ (filled circles) and $B_z = 90 \text{ mG}$ (open circles), and lattice parameters $\Delta = -21\Gamma$ and $U_{pp} = 108E_R$. Lines are the corresponding QMCWF simulations.
CHAPTER 5

RESOLVED-SIDEBAND RAMAN COOLING IN A TWO-DIMENSIONAL OPTICAL LATTICE

5.1 Introduction

This chapter discusses the resolved-sideband Raman cooling experiment which we report in [19], as well as some related experiments which begin with sideband cooled atoms. The basic procedure we use is described in chapter 2. In general, sideband cooling requires resonant Raman coupling. For example, single ions in ion traps have been sideband cooled between ground state hyperfine levels [17, 18]. This requires two "Raman" lasers which are phase-locked but differ in frequency corresponding to the spacing between hyperfine levels, in the gigahertz regime. The beams are frequency tuned to induce stimulated anti-Stokes Raman transitions between the desired levels. A similar method has been applied to neutral atoms in a 1D optical lattice [33]. However, in a 2D optical lattice it is not necessary to sideband cool between hyperfine levels. Experimental complexity can be reduced by sideband cooling between the vibrational levels associated with a pair of Zeeman sublevels. The frequency spacing between vibrational levels is of order 10 kHz. So, if external phase-locked Raman beams are used, they require small relative frequency shifts. More simply, these external beams can be omitted in favor of Raman coupling inherent in the lattice itself [11]. Then magnetic fields can be used to Zeeman shift the desired levels into resonance to enhance the transition rate. This became our favored method because it is both simple and effective.

For sideband cooling, strong Raman coupling is required so that the cooling rate is much larger than the rate of competing heating effects. As noted in chapter 2, Raman coupling between states $|4, m = 4\rangle$ and $|4, m = 2\rangle$ falls off as $1/\Delta_{4,5}$ in a far
off-resonance lattice. Since this is the only Raman coupling which can be induced in a 1D lattice, the prospects for sideband cooling in such a lattice are not good. Indeed we attempted sideband cooling in a 1D lin ⊥ lin optical lattice with little success. Employing both external Raman beams and Zeeman tuning techniques, we were able to induce some limited population transfer between vibrational levels but no significant cooling was achieved [34].

Better results were found in a 2D lattice, where sideband cooling was carried out between states $|4, m = 4\rangle$ and $|4, m = 3\rangle$, for which stronger Raman coupling can be created through the introduction of $\pi$ light. These results alone will be discussed. The 2D lattice, which is made of three beams, is more difficult mechanically to build and align on the MOT than the 1D lin ⊥ lin lattice. Some of the challenges encountered while working with this lattice and their solutions will be detailed. Finally, an ensemble of atoms in a pure state in a far off-resonance optical lattice is a good preparation for initiating experiments that rely upon coherent manipulation, such as momentum squeezing. Since the atoms begin in the same state, they evolve in the same way when the lattice is perturbed. Thus, the signature of coherent evolution is not washed out as it would be if there were several classes of atoms evolving uniquely. The application of sideband cooled atoms to experiments in momentum and position squeezing of the atomic wavepacket is described. Wavepacket oscillations of atoms in optical lattices induced by sudden shifts in potential depth have been used to produce localization below the steady-state value [35], but squeezing below the ground state position spread was not achieved. Since sideband cooled atoms are prepared in the ground state of lattice potential wells, they exhibit the lowest possible momentum variance and kinetic temperature possible for an ensemble of atoms in eigenstates of that particular potential. Use of adiabatic release [36, 29] to reach very cold kinetic temperatures with sideband cooled atoms in this lowest state is also described.
5.2 Technical Considerations

5.2.1 Construction of the 2D Lattice

The three beams which form the 2D lattice are coplanar with the same angle, 120 deg between each pair (figure 2.10), and must intersect at a common point [22]. The lattice plane must contain the gravitational axis so that temperature is measured only in the plane of the lattice in which atoms are trapped. The intersection point of the three beams must be at the MOT position in the vacuum chamber. These conditions can be satisfied by a lattice constructed in the following way. A rectangular alignment frame made from four pieces of precision machined optical rail is assembled so that all four sides meet at right angles and also lie in a plane. Iris apertures, two for each beam, are attached to the frame in such a way that their centers lie in a plane. The position of the irises along the rails is set to define diagonal and vertical beam paths and insure proper intersection of all three. The frame is then centered around the vacuum chamber and MOT and a plumb bob is used to align it with respect to the gravitational axis to about 5 mrad. It is gripped at top and bottom by x, y, z translation stages which permit the frame to be moved and rotated. This allows the beam paths to be centered on the MOT and the lattice frame to be made parallel to the gravitational axis. To produce three lattice beams the far off-resonance laser is split into three beams with combinations of half-wave plates and polarization beam splitter cubes (figure 5.1). The first PBS cube splits the vertical lattice beam from the diagonals while the second evenly splits the remaining power between the diagonal beams. The diagonal lattice beams pass through half-wave plates which allow their polarizations to be rotated into the lattice plane and a series of mirrors is used to align them onto their respective iris pairs on the alignment frame. Polarization beam splitter cubes placed in the diagonal beam paths just before the chamber insure their polarizations lie in a plane. The intensity of the y-axis lattice beam is made ten percent greater than the other beams. A half-wave plate in the vertical beam path rotates the polarization so that the component of polarization in the lattice plane has the same intensity as the diagonal beams. The
Figure 5.1: Polarization optics for producing the 2D lattice beams. Wave-plates allow control of the relative power and polarization of the three beams.

additional power is polarized normal to the lattice plane and is phase shifted by $\pi/2$ with respect to the in-plane component by a quarter-wave plate. So this beam is elliptically polarized. The out-of-plane polarization component introduces $\pi$-light into the lattice which is necessary for the desired Raman coupling between states $|4, m = 4\rangle$ and $|4, m = 3\rangle$ [11]. Mirrors after the polarization optics allow the beam to be aligned on the vertical iris pair. The quantization axis is normal to the lattice plane, taken to be the $z$ axis. The $\sigma^+$ polarized pumper and repumper beams are brought in along this axis from opposite directions to balance radiation pressure. When the pumper and repumper were allowed to propagate in the same direction, unbalanced radiation pressure was found to push the atomic sample over the edge of the probe (temperature measurement) beam. The pumper beam is taken from the master cooling laser and is red shifted by an AOM to make it resonant with the $F = 4$ to $F' = 4$ hyperfine transition. Two magnetic field coils are wrapped around the chamber in the plane of the lattice and are used to produce longitudinal
magnetic fields along the $z$ axis. One coil is used to Zeeman shift optical potentials during sideband cooling and the other is used to preserve the magnetic sublevel distribution of the lattice when the atoms are released and the light is extinguished.

5.2.2 Measuring Temperature in Two Dimensions

The $y$ axis of the lattice is taken to be the lattice beam path into which $\pi$-light is introduced. If the vertical beam path is elliptically polarized, then the $y$ axis of the lattice is along the gravitational axis and we measure the temperature along this axis (figure 5.2). Due to the geometry of the lattice and vacuum chamber, it is not convenient to rotate the lattice by $90^\circ$ deg to measure temperature along the $x$ axis directly. However, it is rather simple to rotate the lattice by $120^\circ$ deg. This is done by introducing elliptical light into one of the diagonal beam paths instead of the vertical path (figure 5.2). Then the $y$ axis is along the diagonal path while the gravitational axis lies along an axis $x'$ that makes an angle of $30^\circ$ deg with respect to the $x$ axis of the lattice. From the temperature along $x'$ the temperature along the $x$ axis is inferred. Due to the isotropy of the wells, the same minimum temperature is expected along the $x$ and $y$ axes and this is what we have observed experimentally.

5.2.3 Alignment and Calibration of Intensity

Determination of whether all atoms are in the lowest vibrational level is made by comparing the kinetic temperature of the sample with the calculated minimum expected temperature, which is called the ground state kinetic temperature, $T_0$. But this calculation depends upon the intensity of the lattice beams and, thus, it is important to measure the lattice intensity very carefully. It is also important that all atoms in the lattice experience the same intensity. This insures that all potential wells will be tuned to the first red sideband and, hence all atoms will be sideband cooled, for the same applied longitudinal magnetic field. Uniformity of intensity is achieved by making the Gaussian profile of the far off-resonance lattice beams much larger than the size of the MOT, which is roughly spherical with a diameter from $300 \mu m$ to $400 \mu m$. The full width at ninety percent of peak intensity for the
Figure 5.2: The 2D lattice with the y-axis along the gravitational axis (a) and rotated by 120 deg (b). Rotation is accomplished by introducing elliptical polarization into a diagonal lattice beam instead of the vertical.
lattice beams is chosen in the range 800 $\mu$m to 1200 $\mu$m. This means that even as the atoms spread out during the sideband cooling, they experience less than 10% intensity fluctuations as long as the MOT sits at the peak of the intensity profile of the three lattice beams. So the lattice beams must be very well centered on the MOT and the on-axis beam intensity must be measured with high accuracy. The procedure is as follows and is pictured in figure 5.3.

A 600 $\mu$m aperture is centered on the far off-resonance laser beam, before it is split into three lattice beams. This insures we use the same part of the beam for alignment of all three paths. Using a power meter to measure the transmitted power, the aperture can be reproducibly superimposed on the beam to better than 100 $\mu$m. The aperture in a lattice beam produces an Airy pattern whose sharp features allow the beam to be well centered on the irises of the alignment frame. Translations of the beam of less than 40 $\mu$m are visible with this technique. A third iris on a translation stage is centered on the beam after the frame and, finally, a 400 $\mu$m aperture is centered on the beam just before the frame. The 400 $\mu$m aperture produces a small Airy pattern at the MOT position. The far off-resonance laser is then tuned to resonance and the position of the lattice beam center becomes evident as atoms are pushed out of the MOT. A mirror after the 400 $\mu$m aperture allows the lattice beam to be adjusted so that it strikes the center of the MOT. The final iris on a translation stage is used to determine the amount of translation of the lattice beam which occurs when adjusting it onto the MOT center and from this the amount of translation of the frame needed to center it on the MOT. Once this is done, the lattice beams need only be centered on the lattice frame itself. Next, the beam intensity is measured at the MOT position. This is done by inserting a mirror into the lattice path to temporarily divert the beam (figure 5.4). A 600 $\mu$m aperture is centered on the beam the same distance from the inserted mirror as the MOT and a power meter measures the on-axis power. Since the power is measured at the MOT position along the beam, any changes in intensity due to mis-collimation are accounted for automatically. The effect of the vacuum chamber window is considered when calculating the intensity. Finally, this alignment and intensity calibration
procedure is applied to the other two lattice beams so that all three are centered on the MOT and have well known intensity. In the end the largest source of uncertainty in the lattice intensity measurement comes from the accuracy of the power meter itself. To insure accuracy, we employed two power meters which had each been recently calibrated by a NIST traceable standard. The meters were tested and found to be consistent in calibration to better than a percent. Before the far off-resonance lattice is split into the three beams, it is combined with near resonance light. Mirrors external to the lattice allow the near resonance beam to be spatially overlapped with the far off-resonance beam. Thus the near resonance lattice is automatically aligned on the MOT and its intensity is measured using the same arrangement as for the far off-resonance lattice. While high accuracy in the intensity calibration is not required for the near resonance lattice, good overlap between it and the far off-resonance lattice is essential for proper transfer between them.

5.3 Running the Experiment

5.3.1 Timing of the Experiment

The sequence of the sideband cooling experiment is as follows. Atoms are trapped and cooled in a 3D MOT and molasses, then transferred into the 2D near resonance lattice for 4 ms where they become localized in the two dimensional potential wells. A longitudinal magnetic field is applied during this time to increase the population in the state \(|4, m = 4\rangle\) in which the sideband cooling occurs [14]. The transfer from near to far off-resonance lattice proceeds just as in the 1D \(\text{lin } \bot \text{ lin}\) optical lattice. The intensity of the near and far off-resonance lattices are linearly turned off and on, respectively, over 500 \(\mu s\). Transfer efficiency is almost 100% and the magnetic sublevel distribution can be well maintained if a relatively large longitudinal magnetic field, of order 1 mG, is applied at the start of the transfer to insure the optical potentials are well separated in energy. Originally, it was thought that a magnetic sublevel distribution in the far off-resonance lattice with all atoms in the state \(|4, m = 4\rangle\) is the optimal preparation for sideband cooling. Cooling on the
Figure 5.3: Experimental setup for centering the lattice frame on the MOT.

Figure 5.4: Technique for lattice beam intensity calibration. The lattice beam is deflected temporarily by an inserted mirror. A 600 µm aperture is placed to simulate the MOT position \( (c = a + b) \) and the on-axis intensity is measured by a power meter.
first red sideband is initiated in the far off-resonance lattice by changing the polarity and magnitude of the longitudinal magnetic field to the "cooling value" which bring states $|4, m = 4, n = 1\rangle$ and $|4, m = 3, n = 0\rangle$ into resonance and by turning on the $\sigma^+$ polarized pumper beam. The $\sigma^+$ polarized repumper beam is on during MOT/molasses, and near and far off-resonance lattice phases. Sideband cooling was attempted for a large number of timings of the longitudinal magnetic field and pumper beam with respect to the near and far off-resonance lattice phases. Best results occur when the longitudinal magnetic field is switched to its cooling value and the pumper beam is turned on about 200 $\mu$s into the transfer, before the far off-resonance lattice has reached its peak intensity. This longitudinal field is not sufficient to maintain the magnetic sublevel distribution of the near resonance lattice upon transfer. Ultimately, sideband cooling was found to be insensitive to this initial distribution. Atoms are sideband cooled in the far off-resonance lattice for 5 to 10 ms. Then, 50 $\mu$s before the FRL light is turned off, the pumper beam is switched off which terminates sideband cooling. At the same time the longitudinal magnetic field is reversed in direction and greatly increased in magnitude to well separate the optical potentials in energy and prevent transitions between Zeeman sublevels while the lattice light is still on. This field also maintains the magnetic sublevel distribution when the lattice light is no longer present. Terminating the cooling well before the lattice is extinguished insures there is no pumper light present when the atoms are released. Any pumper light present after the lattice light is turned off will transfer atoms into the stretched state $|4, m = 4\rangle$ and possibly modify the Zeeman sublevel distribution of the lattice which is one of the properties of the lattice we want to measure. After atoms are released from the lattice their kinetic temperature and Zeeman sublevel distribution are measured as before.

5.3.2 Experimental Parameters

In an attempt to find the optimal conditions for sideband cooling, the experiment was performed for a wide range of parameters. It was found to be sensitive to some of these and relatively insensitive to others. Not surprisingly, sideband cooling is very
sensitive to external magnetic fields. The longitudinal field needed to Zeeman shift the lattice to the 1st red sideband is of order 10 mG. It must be carefully adjusted and maintained to within a few milliGauss to optimize cooling. While not quite as sensitive as the longitudinal fields, transverse magnetic fields must also be well controlled and zeroed. Nonzero transverse fields of only a few tens of milliGauss are large enough to prevent cooling to the ground state. The frequency of the pumper beam is crucial and must be resonant with the \( F = 4 \) to \( F' = 5 \) transition to within a few linewidths. However, the cooling is not very sensitive to its intensity, which can be varied by an order of magnitude with no effect. A typical intensity for the pumper is \( 0.04I_0 \). The repumper beam intensity and frequency were also explored. Cooling is not sensitive to intensity changes of an order of magnitude, but the frequency must be stabilized to a few megahertz. Frequency sensitivity is due to the fact that the potential wells of the \( F = 4 \) and \( F = 3 \) hyperfine ground states are not spatially overlapped. When an atom which is tightly bound in the state \( |4, m = 4) \) optically pumps into the \( F = 3 \) ground state, it is at a saddle point in the two-dimensional potential of this manifold. If it is not repumped back into the \( F = 4 \) state rapidly, its wavepacket begins to disperse causing heating. The frequency of the repumper was set near the \( F = 3 \) to \( F' = 4 \) transition but since this laser is not stabilized to an absolute reference frequency (section 3.2.2), it is prone to drift and must be closely monitored and adjusted often. We also explored the length of the sideband cooling phase and found optimal cooling times are from 5 to 10 ms. Increasing the time beyond this does not improve cooling and causes a slow loss of atoms.

Sideband cooling was carried out at three far off-resonance lattice detunings, \(-2000\Gamma, -3800\Gamma \) and \(-7000\Gamma \), while keeping all other lattice parameters such as polarization, geometry and potential depth fixed. Within the accuracy of the experiment the different detunings have no effect on the cooling, which indicates that for \( \Delta = -3800\Gamma \) the lattice is well into the regime where cooling is not limited by heating due to photon scattering from the lattice. In order to vary the Raman coupling strength the amount of light polarized normal to the lattice plane was
also varied from 5% to 20% of the amount polarized in the lattice plane. Each
time this was changed the in-plane components of polarization in the three lattice
beams were rebalanced. To the level that we can measure, this does not affect the
minimum temperature achieved. This indicates that for our lattice parameters an
out-of-plane component of polarization which has 10% of the intensity of the in-
plane components is sufficient to satisfy the sideband cooling condition of equation
2.29.

5.4 Results of Sideband Cooling

5.4.1 Ground State Kinetic Temperature

The data shown in this section is taken for the following experimental parameters.
The detuning and maximum light shift of the far off-resonance lattice are \( \Delta \approx
-3800 \Gamma \) and \( U_0 \approx 243 E_R \), respectively. The oscillation frequency in the stretched
state \(|4, m = 4\rangle\) is \( \omega_{osc} \approx 20 E_R \). The intensity of the y axis lattice beam polarized
out of the lattice plane is 10% of the in-plane component. This gives rise to Raman
coupling matrix elements between states \(|4, m = 4, n\rangle\) and \(|4, m = 3, n - 1\rangle\) along
the \( x \) and \( y \) axes of about \( 0.7 \sqrt{n_x} E_R \) and \( 2 \sqrt{n_x} E_R \), respectively. The pumper beam
intensity is about \( 0.04 I_0 \) and the length of the sideband cooling phase is about 10 ms.

Equation 2.13 gives the kinetic temperature of a sample of atoms in terms of
the variance of momentum, \( (\Delta p)^2 \), of the sample. In equation 2.16 this is shown
to be expressible as a population weighted sum of the variances of the momentum
probability distributions for the vibrational levels in the well. If all atoms are in
the ground state, this sum reduces to just the variance of momentum of the ground
state, \( (\Delta p_0)^2 \). Since the momentum distribution of the ground state of a harmonic
oscillator is Gaussian, the ground state kinetic temperature, \( T_0 \), is given by,

\[
T_0 = \frac{\Delta p_0^2}{M k_B},
\]

in accordance with equation 2.13. Using equation 2.20 for a harmonic oscillator and
noting that \( \bar{n} = 0 \) when all atoms are in the ground state, \( T_0 \) is given by,

\[
T_0 = \frac{\hbar \omega_{osc}}{2k_B}.
\]  

(5.2)

The oscillation frequencies along the x and y-axes of the potential well corresponding to state \( |4, m = 4\rangle \) are found by independently expanding the potential to second order along x and y. This is similar to expansion of the 1D lin \( \perp \) lin optical potential in section 2.3.3 and the results are also similar. The oscillation frequencies can be expressed in terms of the lattice intensity and detuning. As stated in section 2.4.1, the potential wells are nearly symmetric and the oscillation frequencies are essentially identical along the x and y-axes. Therefore, the ground state kinetic temperature is also the same along both axes. To improve the accuracy of results, anharmonic corrections are applied to \( T_0 \) whenever it is compared to experimental data. [37].

5.4.2 Cooling to the Ground State

Results from sideband cooling are shown in figure 5.5. Plotted are TOF distributions in units of recoil momentum, \( p/\hbar k_L \). The dashed line is for an uncooled sample of atoms held in the far off-resonance lattice for the same duration as for the sideband cooled sample plotted in a solid line. The uncooled sample is prepared by simply withholding the pumper and longitudinal magnetic field for Zeeman shifting to the 1\(^{st} \) red sideband. The sideband cooled sample is an average of TOF distributions for 17 sets of data for cooled atoms taken over the course of about an hour. The circles are not a fit to the data. They represent the calculated ground state momentum distribution for atoms in state \( |4, m = +4, n = 0\rangle \) corresponding to the lattice parameters for which the data was taken. Evidently, the calculation matches the experimental distribution quite well. A Gaussian fit to the experimental data yields a kinetic temperature, \( T_K = 966(10) \text{ nK} \) [38], while the calculated ground state kinetic temperature is \( T_0 = 951(31) \text{ nK} \). Uncertainty in \( T_K \) is obtained from the spread of values of kinetic temperature experimentally extracted from the individual TOF distributions which were averaged together to produce the sideband cooled
Figure 5.5: TOF distributions of atoms in the FRL for an uncooled sample (dashed line) and a sideband cooled sample (solid line). Arrival time has been converted into RMS momentum spread in units of the recoil momentum. The ground state RMS momentum spread calculated for the given lattice parameters is represented by circles.
distribution in the figure as well as the uncertainty in the initial spatial extent of the sample of atoms. The initial spatial extent required a correction of $-45(5)$ nK to the apparent kinetic temperature as discussed in section 3.5.1. The standard deviation in the mean temperature of these sets is about 5 nK as is the uncertainty in the correction due to finite sample size. Uncertainty in $T_0$ is dominated by the accuracy to which the lattice beam intensities can be measured. The intensity is used to calculate the oscillation frequency in equation 2.28.

The experimental and calculated ground state kinetic temperature can be used to find other quantities of interest in the harmonic approximation. The Boltzmann factor is,

$$q_B = \frac{T_K - T_0}{T_K + T_0} = 0.008(16) \quad (5.3)$$

This can be used to calculate the mean vibrational excitation per degree freedom along the $x$ and $y$-axes by

$$\bar{n}_x \approx \bar{n}_y = \frac{q_B}{1 - q_B} \approx 0.008(16) \quad (5.4)$$

The population in the two-dimensional ground state, $\Pi_0$ can also be found from the Boltzmann factor as,

$$\Pi_0 = (1 - q_B)^2 = 0.984(31) \quad (5.5)$$

### 5.4.3 Cooling Time

Another quantity of interest is the time required to sideband cool atoms to the ground state. This is found by changing the sideband cooling time (while holding other parameters fixed) and measuring the temperature at each time. The sideband cooling time is taken to be the time from which the far off-resonance lattice light reaches its peak intensity to the time atoms are released from the lattice. Figure 5.6 shows data taken in this way. Clearly, minimum temperature is reached after about 5 ms and then remains constant. The data is well fit by a decaying exponential with a time constant $\tau = 0.879$ ms.
Figure 5.6: Kinetic temperature as a function of the time for which the sample is sideband cooled is plotted with error bars (circles). The data is fit by a decaying exponential (solid line) of time constant $\tau = 0.879$ ms.

### 5.4.4 Magnetic Sublevel Distribution

The magnetic sublevel distribution of the lattice is measured for sideband cooled atoms by applying a strong magnetic field gradient after the atoms are released from the lattice as described in section 3.5.2. Figure 5.7 shows TOF distributions under Stern-Gerlach analysis for a cooled (upper curve) and uncooled (lower curve) sample of atoms. In the uncooled sample, atoms are relatively evenly distributed over the nine magnetic sublevels, since the magnetic field applied during transfer which appears to optimize sideband cooling does not preserve the magnetic sublevel distribution of the near resonance lattice. For the sideband cooled sample virtually
100% of atoms are in the state $|4, m = 4\rangle$. However, there appears to be some small population in the state $|4, m = 3\rangle$. This is most likely an artifact of the experiment introduced when the sideband cooling is terminated. When the cooling sequence is terminated, the pumper beam is extinguished and a relatively large longitudinal magnetic field, referred to as the alignment bias field, is applied to separate the optical potentials before release from the lattice. As the alignment bias field is applied, potential wells associated with the state $|4, m = 4\rangle$ are moved up in energy with respect to those associated with state $|4, m = 3\rangle$ and vibrational levels between these.
wells become degenerate briefly, which permits atoms to make adiabatic transitions between these states. Evidently, adiabatic transitions can be avoided by switching the magnetic field on so rapidly that vibrational levels do not remain degenerate long enough for transitions to occur or by simply reversing the polarity of the magnetic field so that state \(|4, m = 4\rangle\) is moved down in energy with respect to state \(|4, m = 3\rangle\), thus avoiding any crossing of vibrational level \(|4, m = 4, n = 0\rangle\) with levels in state \(|4, m = 3\rangle\). We were not able to explore the former option due to eddy currents in the UHV chamber and lack of a fast current switching source. The latter is not feasible for the following reasons. During sideband cooling, the magnetic field which tunes the lattice to the 1st red sideband moves state \(|4, m = 4\rangle\) down in energy with respect to state \(|4, m = 3\rangle\). This requires a negative magnetic field. To avoid unwanted vibrational level crossing when the potentials are separated, the applied alignment bias magnetic field must also be negative. Stern-Gerlach analysis is carried out using a magnetic field gradient produced by the MOT coils (see section 3.5.2). The force on an atom undergoing Zeeman state analysis in the experiment is given by equation 3.12. For an atom in state \(|4, m = 4\rangle\), the component of this force along the gravitational (\(z\)) axis is

\[
\mathcal{F}_z = (g_F \mu_B A) \frac{\partial |B|}{\partial z}, \tag{5.6}
\]

This equation holds even if the magnetic field changes direction, as long as it does so slowly enough that the atomic angular momentum can adiabatically follow. The magnitude of the magnetic field produced by the MOT coils increases as atoms fall away from the symmetry axis of the coils, so \(\partial |B|/\partial z > 0\). The sign of the force in this case is such as to accelerate an atom downward into the probe beam which produces an early arriving peak in the TOF distribution. The problem with this is that atoms in states \(|4, m = 4\rangle\) and \(|4, m = 3\rangle\) which are accelerated downward are not well spatially separated at the probe and cannot be properly resolved. There is no problem with TOF resolution if atoms in these states are accelerated upward to produce late arriving peaks. Note that changing the polarity of the magnetic field produced by the MOT coils along the gravitational axis does not help since
\( F_z \) depends only upon the gradient of the magnitude and not the direction of the magnetic field. Therefore, it is not possible to both separate the optical potentials in the desired fashion before release and produce late arriving peaks in the TOF distribution under Stern-Gerlach analysis. The situation is identical for sideband cooling in the state \(|4, m = -4\rangle\).

It should be mentioned that the peak at \( m = 3 \) is not caused by the pumper beam being too weak. Increasing the pumper intensity by more than an order of magnitude for as much as 1 ms in the far off-resonance lattice, even after the potentials have been well separated and the Raman coupling has been shut off, does not eliminate the peak at \( m = 3 \). Finally, it is clear that this peak is not due to mis-alignment of the pumper beam and lattice quantization axis since leaving the pumper on very briefly (40\( \mu \)sec) after the lattice light is extinguished eliminates the peak entirely.

5.4.5 Temperature versus Longitudinal Magnetic Field

Another useful data set is shown in figure 5.8 which plots the sideband cooling temperature as a function of the applied longitudinal magnetic field. As will be shown, this data allows calculation of the ground state kinetic temperature in a way which is independent of lattice beam intensity. The plot reveals the vibrational level structure of the \(|4, m = 4\rangle\) potential as it is Zeeman shifted with respect to the \(|4, m = 3\rangle\) potential and allows corroboration of the lattice intensity measurement. The first minimum in figure 5.8 corresponds to sideband cooling on the 1\(^{st}\) red sideband. Minimum temperature is achieved when states \(|4, m = 4, n = 1\rangle\) and \(|4, m = 3, n = 0\rangle\) are exactly resonant (lower left in figure 5.8). The second minimum corresponds to sideband cooling on the 2\(^{nd}\) red sideband. That is when the states \(|4, m = 4, n = 1\rangle\) and \(|4, m = 3, n = 0\rangle\) are resonant (lower right in figure 5.8). The minimum temperature of this local extremum is not as low as the minimum temperature of the 1\(^{st}\) sideband, since population in the 1\(^{st}\) vibrational level of \(|4, m = 4\rangle\) is not sideband cooled on this resonance. Knowing the spacing in magnetic field between these two minima allows us to calculate the Zeeman shift energy between the 1\(^{st}\) and 2\(^{nd}\) red
Figure 5.8: Sideband cooling temperature versus longitudinal magnetic field is plotted in figure (a). The two local minima correspond to the 1\textsuperscript{st} and 2\textsuperscript{nd} red sidebands. There is an arbitrary offset in the magnetic field. A dashed line indicates the expected minimum temperature. Figure (b) illustrates the two resonances. Vibrational levels of $|4, m = 3\rangle$ are shown as solid lines while those of $|4, m = 4\rangle$ are dashed.
sidebands, which is equivalent to the energy spacing between the first and second vibrational levels of the $|4, m = 4\rangle$ potential. This is given by,

$$|\Delta E_{1,2}| = \frac{1}{4} \mu_B \left| B_{z_1} - B_{z_2} \right|,$$

(5.7)

where $\Delta E_{1,2}$ is the energy spacing between the first and second vibrational levels in the $|4, m = 4\rangle$ potential, and $B_{z_1}$ and $B_{z_2}$ are longitudinal magnetic field values corresponding to the 1st and 2nd red sidebands, respectively. In the harmonic approximation, the oscillation frequency in the $|4, m = 4\rangle$ potential is then given by,

$$\omega_{osc} = \frac{\Delta E_{1,2}}{\hbar},$$

from which the kinetic ground state temperature, $T_0$, is calculated according to equation 5.2. In the anharmonic approximation the same measured $\Delta E_{1,2}$ can be used to calculate the lattice potential depth. Then, $T_0$ can be found from this potential depth in a fully anharmonic model as described in section 2.3.3. Therefore, $\Delta E_{1,2}$ must be determined as accurately as possible. To that end, we introduce a fit model, which involves the Boltzmann factor, to ascertain the magnetic field spacing between the 1st and 2nd red sidebands. Using equation 5.3, we convert each data point in figure 5.8 into a Boltzmann factor and plot $1/q_B$ versus longitudinal magnetic field, which produces two peaks whose maxima are centered on the 1st and 2nd red sidebands as shown in figure 5.9. Note that the Boltzmann factor in equation 5.3 depends upon the ground state kinetic temperature, $T_0$, which is the quantity we want to calculate. However, the particular value of $T_0$ used to calculate $q_B$ initially will not affect the spacing between the two peaks, only their relative amplitudes, and so will not influence the ground state kinetic temperature extracted from this spacing. Thus, it is enough to estimate $T_0$ when calculating $q_B$ for the plot. The fit model used is described as follow. One can show for a harmonic oscillator that the ratio of populations for successively increasing vibrational levels is $1/q_B$. In particular,

$$\frac{\Pi_0}{\Pi_1} = \frac{1}{q_B},$$

(5.8)
where $\Pi_0$ and $\Pi_1$ are the populations of the ground and first vibrational levels, respectively. If we consider only these two levels, then, in steady-state, the principle of detailed balance [39] suggests this ratio is also equal to,

$$\frac{\Pi_0}{\Pi_1} = \frac{\Gamma_{\text{cool}}}{\Gamma_{\text{heat}}},$$

where $\Gamma_{\text{heat}}$ is the heating rate or rate of population transfer from the ground to the first excited state and $\Gamma_{\text{cool}}$ is the cooling rate in the opposite direction. The heating rate is due to scattering of lattice photons, which is independent of the longitudinal magnetic field. The cooling rate is proportional to the strength of the Raman coupling between $|4, m = 4, n = 1\rangle$ and $|4, m = 3, n = 0\rangle$, which we tune with the longitudinal field. Thus, we expect Lorentzian lineshapes about these Raman resonances. Fitting two Lorentzians to the resonances in figure 5.9, allows us to extract the magnetic field spacing between them and, hence, calculate the potential depth and ground state kinetic temperature in a way which is independent of the lattice beam intensity. Using this technique and applying anharmonic corrections, we find $T_0 = 997(50)\, \text{nK}$ which agrees with our calculation based on lattice beam intensity. Given the relative difficulty of measuring intensity, such consistency is reassuring and instills strong confidence in the uncertainty ascribed to $T_0$.

5.4.6 Conclusions

Based on this data we conclude that, to within experimental uncertainty, resolved-sideband Raman cooling has prepared the ground state of an optical lattice. Such a preparation of the lattice represents a pure ensemble of about a million atoms each in separate but identical states both with respect to their internal and external degrees of freedom. The cooling process is found to be sensitive to longitudinal and transverse magnetic fields as well as the frequency of the pumper and repumper beams. However, variation of lattice detuning from 2000$\Gamma$ to 7000$\Gamma$ and of the out-of-plane polarization component from 5% to 20% of the in-plane component has no noticeable effect on the cooling efficiency. The time constant for cooling in our
lattice is a little under 1 ms and we find minimum temperature is achieved by about 5 ms. The result of sideband cooling, that is, the minimum temperature achieved, is found to be robust with respect to both the initial preparation of atoms in the near resonance lattice and the particular characteristics of the transfer of atoms between near and far off-resonance lattices. However, cooling is found to be sensitive to exactly when the longitudinal field for cooling and pumper beam are applied with respect to the transfer. Sideband cooling works so well that it is not necessary to prepare atoms in the near resonance lattice in the stretched state $|4, m = 4\rangle$ with low $\bar{n}$ nor to preserve this preparation during transfer. Ultimately, the experiment was found to work just as well using a constant longitudinal magnetic field set to Zeeman shift to the 1st red sideband in the far off-resonance lattice during both the near and far off-resonance lattice phases. One reason this works is because of an “accidental
degeneracy" of the vibrational levels between which sideband cooling occurs. This means that, for our lattice parameters, states $|4, m = 4, n = 1\rangle$ and $|4, m = 3, n = 0\rangle$ occur at nearly the same energy in the absence of magnetic fields. Therefore, only a small longitudinal magnetic field is required to tune to the first red sideband, and this field can be present during the near resonance lattice phase with little affect on its preparation. This is just one more simplification of the experiment since it does not require dynamically controlled magnetic fields. As expected, rotation of the lattice yielded the same minimum temperatures along the x and y-axes.

An obvious improvement to the experiment is stabilizing the repumper laser frequency to an absolute reference. This would increase the long term stability of the experiment which is crucial if sideband cooling is to be used as a starting point for other experiments.

5.5 Using Atoms Prepared in the Ground State

5.5.1 Momentum Squeezing

One experiment that we attempted with sideband cooled atoms is momentum squeezing. This is done by instantaneously decreasing the potential depth. The change in potential depth is a sudden perturbation and the atoms which are in eigenstates of the initial potential well must be expressed as a superposition of eigenstates in the new well with the expansion coefficients given by the overlap between the original eigenstate and the new ones [35]. The expectation value of the variance of the momentum or position is then time dependent in the new potential well and at certain times can fall below the variance of momentum of the ground state vibrational level in this well, in which case the momentum is said to be squeezed.

The fact that sideband cooled atoms are all in the ground state initially, means that they have the narrowest possible momentum distribution. Since the amount of squeezing of the momentum distribution produced by the change in depth is limited, this improves the chances of finding squeezing below the ground state of the new potential. We measure the variance of momentum directly from the temperature.
The momentum squeezing experiment proceeds as follows. After sideband cooling, the potential depth (lattice intensity) is adiabatically increased to some maximum value (determined by available laser power) and held for about 200 μsec. Then the intensity is instantaneously decreased to a fraction of its original value and the atoms are held at this new depth for a time, t, before being released. Data for the experiment is shown in figure 5.10. Kinetic temperature is plotted against the time, t, for which atoms are held in the shifted or final potential. The data was taken for a sudden shift of lattice intensity from an initial value of 982 $I_0$ to a final value of 172 $I_0$. The ground state kinetic temperature of the final potential is indicated by a line and it is clear that for times around 20 μs the kinetic temperature of the sample is below this value which indicates momentum squeezing for these times. For a given initial intensity, the amount of squeezing should increase with the magnitude of the difference between the initial and final intensities, $I_i$ and $I_f$, respectively. Figure 5.11 shows plots of the maximum momentum squeezing found as a function of the ratio $I_i/I_f$. The amount of squeezing is given by the ratio of the RMS spread of momentum of the sample to the RMS spread of momentum calculated for the ground state of the final potential well. The fact that the amount of squeezing decreases for intensity ratios below 0.2 is an artifact of the experiment caused by atoms being pulled out of very shallow final potential wells by gravity. Experiments in position squeezing were also carried out. The procedure is very similar to momentum squeezing except that the potential depth is suddenly increased rather than decreased. Our experiment is not equipped to measure the position spread, $Δx$, of the atoms directly, but we assume the atoms remain in a minimum uncertainty state. Then $Δx$ can be inferred from $Δp$ using the Heisenberg uncertainty relation. Based upon this assumption we found position squeezing in the shifted potential.

5.5.2 Adiabatic Cooling

Another experiment we attempted with sideband cooled atoms is use of adiabatic release to reach low, sub-recoil kinetic temperatures [36, 29]. After sideband cool-
ing, the lattice intensity or potential depth is adiabatically lowered according to equation 4.1. The theoretical minimum temperature achievable with this technique is determined by the lattice geometry and detuning [29] and can be below the recoil temperature, $T_R$, which is the temperature imparted to an atom upon absorbing or emitting a single photon, about 200 nK.

Experimentally the procedure is complicated by the fact that at very shallow lattice depths atoms are pulled out of the lattice by gravity before the adiabatic release is completed. To counter this effect, the experiment is carried out in a freely falling reference frame. The lattice potential wells are made to accelerate

Figure 5.10: Temperature versus time in the final potential for a sudden lowering of potential depth. The data points (circles), plotted with error bars, are connected by a solid line to illustrate oscillatory behavior. The dashed line indicates the ground state kinetic temperature in the final potential.
Figure 5.11: The ratio of the minimum RMS momentum spread, $\Delta p$, found to the ground state momentum spread of the final potential, $\Delta p_0$, is plotted as a function of the ratio of final to initial lattice intensity. Points for which $\Delta p/\Delta p_0 < 1$ indicate momentum squeezing.
downward along the gravitational axis by varying the frequency of the vertical lattice beam with respect to the diagonal beams. So the diagonal and vertical beams must now be controlled by separate acousto-optic modulators. After the atoms are sideband cooled, the lattice is accelerated and adiabatic release begins. The lattice acceleration is adjusted so that the atoms arrive at the probe beam at the same time as they would if simply released directly after sideband cooling. In principle, the release should take an infinite amount of time, but we truncate it after a period of time, \( t_{\text{rel}} \), called the release time. Adiabatic release was attempted for a variety of values of the release parameter, \( \xi \), (see equation 4.1) ranging from \( 2 \text{ms}^{-1} \) to \( 50 \text{ms}^{-1} \) and release times ranging from \( 0.1 \text{ms} \) to \( 10 \text{ms} \). Best results were found for \( \xi = 10 \text{ms}^{-1} \) and \( t_{\text{rel}} = 2 \text{ms} \). Figure 5.12 shows TOF distributions for sideband cooled atoms and for sideband cooled atoms with adiabatic release. After compensating for the finite size of the MOT, the minimum adiabatic release temperature found was \( T_{\text{min}} \approx 120 \nK \).

In the harmonic model, the ground state kinetic temperature is proportional to the oscillation frequency (equation 5.2) and, therefore, the square-root of the lattice potential depth (equation 2.28). The minimum temperature under adiabatic release in the harmonic model asymptotically approaches zero. However, in an optical lattice as the uncertainty in \( x \) approaches the spacing between optical potential wells, the harmonic oscillator approximation is no longer valid and band theory must be applied to calculate the final kinetic temperature under adiabatic release. This is done by expressing the initially tightly bound harmonic oscillator states in the basis of tight-binding Bloch states and allowing these to be adiabatically transformed into free Bloch states [29, 10]. Since this is done adiabatically, the population in each Bloch state is preserved. The final temperature is assigned by making a population weighted average of \( p^2/2M \) over the momentum distribution for the free particle Bloch states and equating this to \( k_B T_f/2 \), where \( T_f \) is the final temperature achieved at time \( t \to \infty \). The final temperature depends upon the populations of the Bloch levels which are found from \( \bar{n} \) for the initial sample and the lattice periodicity. For an initial thermal distribution with \( \bar{n} = q_B = 0 \), then \( T_f = T_{\text{min}} \), where \( T_{\text{min}} \) is
Figure 5.12: Momentum distributions in units of the recoil momentum for atoms which are sideband cooled only (dashed line) and for atoms which are adiabatically released from the lattice after sideband cooling (solid line).
the minimum kinetic temperature possible under adiabatic cooling in the lattice. For our two-dimensional lattice \( T_{\text{min}} \approx 41 \text{nK} \) [40]. Since for sideband cooled atoms \( \bar{n} \approx 0 \), we hoped to achieve this minimum kinetic temperature.

It is unclear what prevents the experiment from reaching the theoretical minimum temperature. Experimental artifacts which might broaden the measured temperature such as the MOT size or probe beam thickness were reconsidered. The shape and smoothness of the adiabatic release waveform was carefully checked with much attention paid to the low intensity region in the extended tail of the release and also to how the waveform was truncated. However, the minimum temperature was found to be relatively insensitive to changes in this waveform, and \( \xi \) and \( t_{\text{rel}} \) could be varied over a wide range with little affect. Properties of the lattice were also considered. The intensity among the lattice beams was misbalanced by a few tens of percent with neutral density filters. Also, frequency fluctuations at least ten times larger than could be expected from the laser itself were imposed on the lattice light during adiabatic release. Neither of these rather large perturbations to the lattice had any significant effect. Numerical simulations of adiabatic release suggest that large frequency and amplitude fluctuations of the lattice light as well as oscillation of the spatial position of the potential wells by significant amounts (of order \( \lambda \)) have only a small effect on the minimum temperature [40]. It seems unlikely, therefore, that the minimum temperature we achieve is limited by laser noise or by shaking of the lattice due to mechanical vibrations of mirrors. Recently, we discovered that retroreflections off optical elements in a lattice beam path such as a vacuum chamber window can create standing waves that interfere with the motion of atoms in an accelerating lattice. In the two dimensional lattice, for example, a reflection off an optical element in the vertical beam path could produce a standing wave superimposed on top of the accelerating lattice. In principle, the standing wave could delay some of the atoms which would increase the spread in arrival times of all atoms at the probe beam and be perceived as an increase in kinetic temperature.

Experimental data for adiabatic release is compared to the harmonic model in figure 5.13. The upper plot shows the temperature as a function of the time into
the release for experimental data (circles) and for the harmonic model (line). In both cases the initial temperature and adiabatic release parameter are $T_i = 976 \text{nK}$ and $\xi = 20 \text{ms}^{-1}$, respectively. A dashed line indicates the expected experimental minimum temperature of $41 \text{nK}$. The data begins to exhibit noticeable departure from the harmonic model after about $0.5 \text{ms}$ and reaches a minimum temperature of around $130 \text{nK}$ in this case. The same data points are plotted against the square-root of the ratio of the final to initial potential depths in the lower plot. The harmonic model is indicated by the straight line and departure of the experimental data (circles) from this becomes evident when the potential depth falls to a little under $20\%$ of its initial value.
Figure 5.13: Temperature data (circles) is plotted versus (a) the time into adiabatic release and (b) the square root of the ratio of final to initial potential depth for adiabatic release parameter $\xi = 20 \text{ms}^{-1}$. The solid line indicates the expected behavior for adiabatic lowering of the ground state of a harmonic oscillator, and the dashed line indicates the expected minimum temperature calculated from a band theory model.
CHAPTER 6

CONCLUSION

Optical lattices provide a rich system for studying the atom-light interaction in a regime where the field intensity is low, and the kinetic energy of the atom is comparable to the light shift interaction energy between atom and field. An atom tightly bound in a potential well of a far off-resonance lattice is a model quantum system for which dissipation can be largely suppressed. This offers the hope of exerting great control over an atom's motional and internal states. Optical lattices are highly flexible. Their potential wells can be varied in dimension, shape, spacing and depth by choice of laser beam geometry, polarization, intensity and detuning, and by the application of magnetic fields. Coupling between atomic internal states can also be introduced in a selective way by choice of laser beam polarization and magnetic field. Freedom in the design properties of a lattice can aid in the preparation and control of atoms in the lattice.

A wide range of experiments in lattices would benefit from starting with atoms prepared in the vibrational ground state of the lattice potential wells. These include momentum and position squeezing, the generation of Fock and coherent states and tunneling between double well potentials in a lattice [11]. Also, the lowest kinetic temperature that can be reached with adiabatic cooling in an optical lattice is achieved when atoms in the lattice are prepared in the vibrational ground state.

The primary experiment described in this dissertation is resolved-sideband Raman cooling in a two-dimensional optical lattice. Experiments in a one-dimensional lin ⊥ lin optical lattice which are relevant to sideband cooling were also described. These auxiliary experiments include transferring atoms from near into far off-resonance lattices, determining heating rates in a far off-resonance lattice, and improving laser cooling in a near resonance lattice by the application of weak magnetic fields. The
importance of these experiments to sideband cooling and their results can be summarized as follows.

To be effective, sideband cooling must be carried out in a far off-resonance lattice with low initial mean vibrational excitation, $\bar{n}$. Since built-in laser cooling is not present in a far off-resonance lattice, we have attempted to transfer atoms between near and far off-resonance lattices, while maintaining the low $\bar{n}$, characteristic of a near resonance lattice. The desired result was achieved by using a transfer method that involves maintaining tight binding of atoms as they go from near into far off-resonance potential wells. This is accomplished by overlapping NRL and FRL potential wells and linearly decreasing the near resonance lattice intensity, while increasing the intensity of the far off-resonance lattice. This technique causes negligible heating and roughly preserves the magnetic sublevel distribution of the near resonance lattice.

The minimum temperature which can be reached by sideband cooling represents a balance between the sideband cooling rate and the heating rate due to photon scattering. Cooling to the ground state, $\bar{n} = 0$, requires that the cooling rate be much greater than the heating rate. The heating rate in a 1D lin $\perp$ lin far off-resonance optical lattice was measured and found to be much less than the anticipated sideband cooling rate. This increased our confidence in the prospects of sideband cooling to the vibrational ground state. Of secondary importance, the heating rate was found to be nonlinear at long times in the far off-resonance lattice. Although we were not able to verify it, the data suggests that $\bar{n}$ may come to steady-state, which would indicate the presence of a cooling mechanism in the FRL [30].

Originally, it was thought that the results of sideband cooling in the far off-resonance lattice would be sensitive to how the atomic sample is precooled in the near resonance lattice. We attempted to produce a sample with a high population in the ground vibrational level of a single Zeeman stretched state, by application of weak longitudinal and transverse magnetic fields. Longitudinal fields were found to increase this population by 33% over the zero field value. Furthermore, for a given longitudinal field, transverse fields were found to significantly reduce $\bar{n}$. Although
it was not measured directly, simulations suggest that a judicious application of transverse and longitudinal fields can increase the ground state population in a single Zeeman sublevel by 80%. Ultimately, sideband cooling results were found to be relatively insensitive to pre-cooling in the NRL. However, improving laser cooling in the NRL with weak magnetic fields has already proven useful in experiments that produce a pure state through selection of the ground state of a lattice.

Resolved-sideband Raman cooling has been used to prepare atoms in the vibrational ground state of a two-dimensional optical lattice. When prepared in this fashion, atoms in the lattice are in identical, pure quantum states both with respect to their motional and internal degrees of freedom. Sideband cooling occurs between vibrational levels in Zeeman substates associated with a single hyperfine level, and Raman coupling is provided by the lattice beams themselves. A longitudinal magnetic field tunes the lattice to the first red sideband to initiate cooling. We found sideband cooling prepares 98% of atoms in the two-dimensional ground state with mean vibrational excitations along the x and y-axes of $\bar{n}_x \approx \bar{n}_y \approx 0.008(16)$ and corresponding Boltzmann factor $q_B = 0.008(16)$. Zeeman sublevel analysis indicated that virtually all atoms are in the stretched state $|4, m = 4\rangle$ as well. The ground state kinetic temperature, $T_0$, which is used to determine $\bar{n}$ and $q_B$, was calculated by two independent methods. This instills confidence in the uncertainties given for these quantities. Overall, results of sideband cooling were found to be remarkably insensitive to many experimental parameters such as lattice detuning, the amount of Raman coupling present, and the intensity of the pumper beam. Results were also insensitive to pre-cooling of atoms in the NRL, so it was possible to simplify the experiment by using only static magnetic fields.

Two experiments were carried out using sideband cooled atoms, momentum squeezing and adiabatic cooling. Using momentum squeezing, the RMS momentum spread of the sample was reduced below the spread of the ground state by about 10%. Adiabatic cooling, starting with all atoms in the vibrational ground state, was carried out in a freely falling reference frame, and resulted in a minimum temperature of around 120 nK, about 3 times the expected minimum value. We have no firm
evidence to explain this disparity, but suspect the presence of a stationary lattice created by uncontrolled reflections from optical components in the lattice beam path.

Moreover, sideband cooling is important because it has, for the first time, prepared atoms in an optical lattice in a pure quantum state. As mentioned, preparing atoms in the vibrational ground state of lattice potential wells may be useful in experiments involving coherent control, such as squeezing and tunneling. Future work will include extending this sideband cooling scheme into a 3D lattice. More practically, atoms cooled by adiabatic expansion after sideband cooling in 2 or 3 dimensions, may find use in the construction of improved atomic clocks.
APPENDIX A

EXTERNAL CAVITY LASER DESIGN

The master laser for the high power laser system is based upon an external cavity laser suggested in a paper by Wieman [24] but includes significant improvements in mechanical design. The modified laser (figure A.1) combines off-the-shelf mechanical components with a simple copper mount to make a laser which is small, easy to build and, more importantly, has considerably better long term mechanical stability than its predecessors. The laser is built around the Thorlabs LT230A collimation tube assembly which is a precision machined mount made to rigidly hold a standard 9 mm diode laser package and an aspheric collimation lens. The diode and lens are aligned radially by the tube while the lens, which is in a threaded mount, can be screwed in or out to collimate the beam. An O-ring compresses the lens mount to maintain collimation once it is set. A simple copper mount with a through hole firmly grips the collimation tube and provides a heat sink for the laser. One face of the copper mount is machined at 30 deg to match the angle of the first diffracted order from the grating. To this face is applied a stainless steel flexure mount (Newport MFM100) with tip and tilt motion. A steel strut is connected to the flexure mount and the piezoelectric transducer and diffraction grating are epoxied to the strut.

Overall, the copper mount holds the collimation tube at the correct angle with respect to the grating for Littrow configuration. The flexure mount allows horizontal and vertical fine tuning of the feedback while providing excellent long term pointing stability. As before, the PZT controls the cavity length. The copper mount rests atop a 1" square thermoelectric cooler for temperature stability.

Moreover, the advantages of this design are many. The modular construction allows collimation before the tube is even integrated into the mount. Use of off-the-shelf parts reduces cost and machining time. Compact size allows the laser
to come to thermal equilibrium rapidly when the laser is turned on. The rigid construction maintains the feedback from the grating so consistently that the laser can be operated for well over six months without adjustment as compared to a few weeks with our earlier designs.
Figure A.1: Design of modified external cavity laser. All lengths given in inches.
Assembly Instructions

1) Epoxy piezoelectric transducer to stainless steel strut.
2) Epoxy grating to piezoelectric transducer.
3) Attach stainless steel strut to flexure mount.
4) Attach stainless steel flexure mount to copper mount.
5) Collimate laser diode.
6) Insert collimation tube assembly into copper mount.

Figure A.2: Assembly instructions for external cavity laser.
Figure A.3: External cavity laser after assembly. The direction of the output beam is indicated.
REFERENCES


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