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The force on an atom moving in an intense standing wave

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THE FORCE ON AN ATOM MOVING IN AN INTENSE STANDING WAVE

by

JEFFREY JOHN TOLLETT

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APPROVED, THESIS COMMITTEE

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Abstract

Atoms moving in an intense, near resonant standing wave (SW) experience a force along the direction of the SW. This force has been investigated by observing its effect on the transverse velocities of atoms in a thermal beam. When the standing wave is tuned above resonance the atomic beam is collimated and has an increased intensity at the center of the beam. Below resonance, the SW anti-collimates the atomic beam, reducing the central intensity. These effects are due to the interaction between the induced atomic dipole and the inhomogeneous field of the SW. Further investigation shows structure in the transverse velocity distribution of the atomic beam. This structure is due to multiphoton (Doppleron) resonances in the force. Calculations reproduce the dipole force and multiphoton resonance effects and are in good agreement with the data.
Acknowledgements

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I. Introduction

Radiation pressure results from the force exerted on an atom by an electromagnetic field. Given an appropriately monochromatic, tunable source, this force can be used to reduce the velocity of an atom, and therefore reduce its temperature. Using such a process it is possible to cool a gas of atoms to well below 1 K. The reduction of Doppler and transit time broadening effects greatly benefit the study of atomic structure. In addition, the large deBroglie wavelength, $\lambda_b$, of cold atoms is expected to be important in collisions between cold atoms. An effect which has already been seen is a dramatic increase of the cross-section for a certain inelastic collision process.\(^1\) If $\lambda_b$ is comparable to the interatomic separation and the atoms are sufficiently non-interacting, one may observe collective effects predicted by quantum statistical mechanics, such as Bose-Einstein condensation.\(^2\)

Radiation pressure can also be used as a tool to collimate, bend or focus atomic beams. Transverse cooling of an atomic beam can be used to collimate the beam while increasing the on-axis beam intensity. Similarly, there are several ways one can redirect an atomic beam. Bending angles as large as thirty degrees have been seen for a cold beam,\(^3\) while in our laboratory angles of up to three degrees have been observed for a thermal beam.\(^4\) It is possible with the
use of radiation pressure techniques to focus an atomic beam to a point\textsuperscript{5} or funnel a large beam into a narrow high intensity beam.\textsuperscript{6} Others have described manipulating macroscopic bodies using radiation pressure.\textsuperscript{7}

The classical basis for radiation pressure is presented by Maxwell in his *Treatise on Electricity and Magnetism,* where he shows that when an electromagnetic wave is incident upon a surface the stress tensor evaluated on the surface gives a pressure, $P$, which depends on the intensity of the radiation, $I$, the reflectivity of the surface, $r$, and the speed of light $c$;

$$P = \frac{I(1+r)}{c}$$

This is equivalent to the conservation of momentum density at the surface. The relationship is known as the Maxwell-Bartoli formula since Bartoli independently predicted the phenomenon.

Owing to the great difficulty in the measurement of this pressure (4.7$\times$10$^{-6}$ N/m$^2$ for sunlight on earth) it wasn't until the turn of the century that Lebedev\textsuperscript{9} and then Nicols and Hull\textsuperscript{10} finally experimentally demonstrated agreement with the theory with sufficient accuracy.

The next step was taken by Einstein who hypothesized the concept of photons. As the particle-like description of EM radiation, photons carry momentum.\textsuperscript{11} Photons with energies corresponding to allowed transition energies of an atom will
be resonantly scattered by stimulated absorption and stimulated and spontaneous emission. Each event involves a momentum exchange between the atom and photon. Stimulated events transfer momentum along the direction of the incident radiation. However, spontaneous emission is isotropic when averaged over many events. Consequently, the momentum transferred to the atom by spontaneous emission averages to zero. If all the absorbed photons come from the same direction, it is evident that there will be a net impulse supplied to the atoms in the direction of the incident radiation.

Otto Frisch preformed an atomic beam experiment in 1933 in which he measured the deflection of an atomic sodium beam due to the effects of a sodium resonance lamp.\textsuperscript{12} In this case, he set out to measure something significantly different from the previous radiation pressure experiments, namely, resonance radiation pressure (RRP). One would expect a considerable enhancement in the magnitude of the force for a given intensity of radiation if the light is in resonance with a particular atomic species.

Interest in the use of resonance radiation pressure (RRP) grew with the availability of high intensity, narrow bandwidth, tunable light sources (lasers). In 1975, the concept of laser cooling of gases was invented simultaneously by Wineland & Dehmelt and Schawlow & Hängsch\textsuperscript{13}.

This thesis concentrates on the concept of cooling atoms
along one axis. In the case of an atomic beam, if the cooling were along the direction transverse to the beam axis, this would correspond to collimation. Consider the action of a low intensity, near resonant laser with frequency \( \omega_L \), propagating transverse to an uncollimated atomic beam. Atoms with resonance frequency \( \omega_0 \) moving with some transverse velocity, \( v_x \), will see the Doppler shifted frequency, \( \omega \),

\[
\omega = \omega_L - k \cdot v_x
\]

where \( k \) is the laser propagation vector. For light tuned below resonance atoms counter-propagating to the laser will be Doppler shifted into resonance. As the atoms are cooled they are Doppler shifted out of resonance. The average transverse velocity is reduced and these atoms have been cooled. It is evident that atoms co-propagating with the laser will experience a heating force due to the finite excitation linewidth. The force on the co-propagating atoms is smaller than that for counter-propagating atoms because of the larger detuning. Simultaneous cooling of both positive and negative velocity wings of the transverse distribution can be achieved by retro-reflecting the laser. For low intensity the forces of the opposing beams add giving a force directed towards \( v_x = 0 \), where the force is zero. This arrangement is called "optical molasses". The reason being that RRP is not a spatial restoring force, it is a velocity dependent damping force. Atoms in the field of the opposing laser beams are prevented from moving along either beam by RRP.
The low temperature limit, for a two-level system, is determined by the random emission process giving,\(^{15}\)

\[ T_{\text{min}} = \frac{\hbar \gamma}{2k} \]

where \( \gamma \) is the spontaneous decay rate of the upper level. By use of the equipartition theorem one can see that the RMS transverse velocity, \( v_{\text{rms}} \), and temperature, \( T \), are related by,

\[ v_{\text{rms}} = \sqrt{\frac{kT}{m}} \]

This gives a minimum RMS velocity of,

\[ v_{\text{rms}}^{\text{min}} = \sqrt{\frac{\hbar \gamma}{2m}} \]

For the lithium D line (2S-2P), \( T_{\text{min}} = 141 \, \mu K \), and \( v_{\text{rms}}^{\text{min}} = 0.40 \, \text{m/s} \). The magnitude of the RRP can be found as follows. A two-level atom will exchange at most a momentum of \( \hbar k \) with a single photon. The fastest this can occur is limited by the spontaneous decay rate, \( \gamma \). There is an addition factor of \( 1/2 \) since this is the maximum excited state population. Therefore,

\[ F_{\text{max}} = \frac{\hbar k \gamma}{2} \]
Physically, the maximum momentum exchanged is determined by the excitation frequency, while the maximum rate is limited by the spontaneous decay rate (stimulated absorption followed by stimulated emission yields no net momentum exchange).

When the atoms are placed in an intense near resonant radiation field a new description must be used to explain the behavior. If the intensity is increased beyond saturation, the two atomic states are not eigenstates of the total Hamiltonian since the level energies are shifted due to the AC Stark effect. In the case of a standing wave, the level shifts vary spatially along the standing wave giving rise to a force due to the interaction between the induced atomic electric dipole moment and the inhomogeneous electric field. This is known as the gradient or dipole force.\(^6\) This force can be much larger than RRP for small velocities reducing the required time to cool an atom.

For certain velocities it is possible to make the transition to the excited state via a multiphoton process involving stimulated emission and absorption by both travelling wave components of the standing wave. The effect of this process is to give sharp resonances in the force at particular velocities. These multiphoton resonances involve the transfer of energy quanta, \(kv\), dubbed the "Doppleron", and therefore are often referred to as Doppleron resonances.\(^7\) The study of these features involves more careful study of evolution of the atom in the perturbing potential.
II. Theory of Laser Cooling in a Standing Wave

The solution of the equations of motion of a two-level atom moving in a standing wave can be found by means of the density matrix. These solutions are exact and calculations of the force yield structure in the force which depend on the intensity and detuning of the standing wave. Unfortunately, this method gives no physical picture of the interaction. The results can be interpreted in three regimes: 1) the dipole force, 2) multi-photon or Doppleron regime and 3) the usual RRP regime. These effects can clearly be seen in calculations of the force based on the density matrix solutions. The theory of dressed states as put forth by Cohen-Tannoudji and Dalibard offers the best physical picture of the dipole force on the atom. While the picture of Kyrölä and Stenholm offers the best physical picture of the Dopplerons. First an outline of the general solution in terms of the density matrix will be given. A brief description of the other pictures will then be given with the corresponding physical insights.

II. A. Solution of the Equations of Motion (Density Matrix)

The density matrix is a representation of the expectation of the state of the system, in a given basis, at any point in time. It is a square, Hermitian matrix which is N x N, where
N is the number of states which span the space. Diagonal terms, \( \rho_{kk} \), refer to the populations of the states while the off-diagonal terms, \( \rho_{nk} \), indicate the coupling between states. When the basis states are eigenstates of the unperturbed Hamiltonian \( (H_0|n> = \hbar \omega_n|n> \) ) the Liouville equation with phenomenological damping leads to the equations of motion for the matrix elements,

\[
\dot{\rho}_{km} = \rho_{km}(-i\omega_{km}-(\gamma_k+\gamma_m)/2)-\frac{i}{\hbar} \sum_n (V_{kn}\rho_{nm}-\rho_{kn}V_{nm})
\]

\[
\dot{\rho}_{kk} = -\gamma_k\rho_{kk} + \sum_n \gamma_{nk}\rho_{nn} - \frac{i}{\hbar} \sum_n (V_{kn}\rho_{nk}-\rho_{kn}V_{nk})
\]

where \( \gamma_k \) is the total decay rate from the \( k^{th} \) level, \( \gamma_{nk} \) is the decay rate from the \( n \) to the \( k^{th} \) level, \( \omega_{km} = \omega_k - \omega_n \) and \( V_{kn} \) is the coupling between levels \( k \) and \( n \). The dot notation refers to a total time derivative,

\[
\frac{d}{dt} = \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla
\]

which shows how the velocity dependence is incorporated into the density matrix. We are interested in a two level system, therefore \( k, n \) and \( m \) are either elements of \( |e> \) or \( |g> \), the excited or ground states. We also make the dipole approximation such that the interaction potential is simply,

\[
\mathbf{V} = -\mathbf{d} \cdot \mathbf{E}
\]
For a standing wave the electric field can be written as,

\[ E = 2E_0 \cos(kz) \cos(\omega_L t) \]

where \( E_0 \) is the magnitude of each travelling wave field. This yields, in the rotating wave approximation,

\[
\dot{\rho}_{ee} = 2 \Omega \cos(kz) \text{Im}(\sigma_{eg}) - \gamma \rho_{ee} \\
\dot{\sigma}_{eg} = (i\Delta - \frac{\gamma}{2}) \sigma_{eg} + i\Omega \cos(kz) (1 - 2\rho_{ee})
\]

where, \( \Omega = 2dE_0/\hbar \) is the Rabi frequency, \( \Delta = \omega_L - \omega_e \) is the detuning and \( \sigma_{eg} = \rho_{eg} \exp(-i\omega_L t) \). This approximation is equivalent to a time average over a period, \( 1/\omega_L \) and ignores terms of frequencies \( 2\omega_L \). Similarly the time average force on the atom over one period is,

\[
f(x) = \langle -\nabla_x V \rangle \\
= \nabla_x \langle E \rangle \langle d \rangle \\
= \nabla_x \langle E \rangle \text{Tr}(\rho d) \\
= -\hbar \Omega \sin(kz) \text{Re}(\sigma_{eg}) \hat{Z}
\]

The force can be found by integration of the equations of motion of the density matrix or by the method used by Stenholm and Minogin.\(^9\) In their method the above set of equations are re-written in terms of the Bloch type variables,
\[ r = 1 - 2 \rho_{ee},\]
\[ s = 2 \text{Im}(a_{eg}),\]
\[ c = 2 \text{Re}(a_{eg}).\]

This gives,
\[ t = \gamma (1-r) - 2 \Omega \cos(kz) s,\]
\[ \dot{s} = \Delta c - \frac{\gamma}{2} s + 2 \Omega \cos(kz) r,\]
\[ \dot{c} = -\frac{\gamma}{2} c - \Delta s.\]

One then looks for steady state solutions which are periodic with wavelength,
\[ r(v,z) = \sum_{n=-\infty}^{\infty} r_n(v) e^{inkz} \]

with similar expressions for \( s(v,z) \) and \( c(v,z) \). When these are substituted back into the above equations one gets recurrence relations between the various coefficients \( r_n(v) \), \( c_n(v) \) and \( s_n(v) \). These recursion relations are then solved in terms of a continued fraction. The continued fraction gives a rapidly converging evaluation of the steady state populations which depend on \( \Omega, \Delta, \gamma, k \) and, \( v \). This solution is identical to directly integrating the density matrix equations of motion to find the periodic variation of the
population. The main benefit of the continued fraction solution is that the computation time can be several orders of magnitude faster than directly integrating the equations of motion.

Calculations of the force are given in figure 1, which show the development of the force for various choices of the Rabi frequency, $\Omega$, and the detuning, $\Delta$, as a function of the relative velocity, $v$, between the atom and the standing wave. For a given $\Delta$, it is clear that as $\Omega$, which is proportional to the field strength, is increased the force evolves from the usual RRP (seen at low intensities) to something very different (figure 1(a)). For the small values of $\Omega$, the standing wave force is the sum of RRP from the opposing travelling waves. This is the small positive peak in figure 1(a) for $\Omega = \gamma$. For larger value of $\Omega$, a large, narrow peak opposite in sign of the RRP develops at small $v$. This feature is the dipole force. Structure due to multiphoton processes also begins to occur, manifesting itself at higher velocities than the dipole force peak. The locations of these peaks shift to larger $v$ as $\Omega$ increases. This is due to Stark shifts of the energy levels. They also broaden due to power broadening effects. These points are important in the experiment and will be elaborated on later. For positive $v$ and negative $\Delta$ the dipole force is opposite in sign from the RRP. The sign of the force shown in figure 1 reverses upon reversal of the sign of $\Delta$ or $v$. 
Figure 1 The variation of the force on an atom in a standing wave in units of maximum RRP, as a function of the relative velocity, $v$, between the atom and the standing wave. (a) shows the development of the Doppler resonances and dipole force with increasing $\Omega$ ($\propto$ field strength) at a given $\Delta$ ($=\omega - \omega_0$). As the field strength increases structure appears in the force. (b) shows that $\Delta$ changes the position and magnitude of the structure.
II. B. Dipole Force

An atom placed in an electric field develops a dipole moment. If this field has a gradient at the atom then the atom will see a force, \( f = \nabla (d \cdot E) \). The dressed atom picture offers a physical understanding of this force when the atom is moving in a standing wave. The dressed atom refers to a new set of eigenstates which come from diagonalizing the total Hamiltonian, including the internal energy of the atom, the field mode energy, and the atom field coupling energy. The new energy eigenstates depend on the number of photons in the field, \( n \), as well as \( \Delta \) and \( \Omega \). Figure 2(a) shows how a typical set of these eigenstates vary with \( \Omega \) for \( \Delta > 0 \). The dressed basis can be expressed in terms of the basis \((|e>, |g>, n)\) as,

\[
|1, n\rangle = \cos \theta |e, n\rangle + \sin \theta |g, n+1\rangle \\
|2, n\rangle = -\sin \theta |e, n\rangle + \cos \theta |g, n+1\rangle
\]

with,

\[
\cos 2\theta = -\frac{\Delta}{\Omega'}, \quad \sin 2\theta = \frac{\Omega}{\Omega'}, \quad \Omega' = \sqrt{\Omega^2 + \Delta^2}
\]

and energy levels,

\[
E_{1n} = (n+1) \hbar \omega_L - \frac{\hbar \Delta}{2} + \frac{\hbar \Omega'}{2} \\
E_{2n} = (n+1) \hbar \omega_L - \frac{\hbar \Delta}{2} - \frac{\hbar \Omega'}{2}
\]  \quad (1)
Figure 2  (a) The dressed states connect with the unperturbed ground states in the absence of a field. For $\Delta > 0$ the state $|1, n\rangle \rightarrow |g, n+1\rangle$ and $|2, n\rangle \rightarrow |e, n\rangle$. As the field increases the dressed states become mixtures of $|g, n+1\rangle$ and $|e, n\rangle$. Note that $\omega_o \gg \Delta$.  
(b) For $\Delta > 0$ atoms emit preferentially from the peaks of the potential hills resulting in a net cooling force.
The main point to follow here is that out of the field the state \(|1, n>\) (\(|2, n>\)) correlates to \(|g, n+1>\) (\(|e, n>\)) for positive \(\Delta\) while just the opposite is true for negative \(\Delta\). In other words, the \(|1, n>\) state is more contaminated by the \(|g, n+1>\) state for positive detuning than for negative detuning. This condition is reversed for the opposite sign of detuning.

The net force on this atom can be written as the sum of the forces in the states \(|1, n>\) and \(|2, n>\) weighted by the populations in these states, \(\rho_{11}\) and \(\rho_{22}\) respectively:

\[
f = -\rho_{11} \nabla E_{1n} - \rho_{22} \nabla E_{2n}
\]

From equation 1 one can see that,

\[
\nabla E_{1n} = -\nabla E_{2n}
\]

so that,

\[
f = \nabla E_{1n} (\rho_{22} - \rho_{11})
\]

For \(\Delta > 0\) we have \(\rho_{11} > \rho_{22}\) since \(|1, n>\) is mostly ground state. Therefore, the static force repels the atoms from the standing wave field anti-nodes. A similar argument applies for \(\Delta < 0\) such that the atom is attracted to the field anti-nodes.

As the atom moves through the standing wave it samples various instantaneous forces. The effective force the atom feels is the average of the instantaneous forces. For \(\Delta > 0\) the atom feels a negative force. This can be understood by
following the atoms through several spontaneous emission cycles (figure 2(b)). An atom in the state $|1,n\rangle$ will most likely radiate at a field antinode where it is most contaminated by $|e\rangle$ (the undressed excited state). The spontaneous emission will most likely place the atom in the $|2,n-1\rangle$ state, where it is most likely to radiate at a field node, and so forth. In each case the atom climbs a potential hill and radiates from the top falling into a valley. On average the atom will be slowed by moving up more potential hills than it comes down. Likewise when $\Delta < 0$ the atom is heated.

The major advantage of the dipole force is that it does not saturate with intensity like RRP, instead the force is proportional to $\Omega$ and therefore, increases as the square root of the field intensity. This explains why the dipole force peaks in figure 1 can be much larger than RRP.

II. C. Multiphoton Resonances

An atom moving in a standing wave will see the Doppler shifted frequencies of the standing wave, $\omega^\pm = \omega_\pm \pm kv_x$, corresponding to the two different travelling wave components of the standing wave. The simplest absorption process occurs when $|\Delta| = kv$. This is the case of one photon resonant absorption. However, more complicated processes are possible. The transition from the ground state to the excited state can
occur by means of multiphoton processes. Figure 3(a) shows this process for $\Delta < 0$. The atom can absorb a photon from one travelling wave to a virtual level, be stimulated to emit a photon by the second travelling wave and absorb another photon from the first travelling wave. This three photon process is shown in fig. 3(b). The five photon process is shown in fig. 3(c). The doppler shifted frequencies of the travelling wave components of the SW are,

$$\omega^*_n = \omega_L \pm kv_n$$

When the net change in energy from the $n+1$ absorptions and $n$ emissions equals the energy level separation of the excited and ground state, $\Delta \omega_e$, the process is resonant. The approximate velocity at which a $2n+1$ order process is resonant is given by,$^{17}$

$$kv_n = \frac{-\Delta}{2n+1} \quad (2)$$

The effect is that the detuning is compensated for by the absorption or emission of $2n+1$ units of $kv$, or $2n+1$ Dopplerons. This form is valid only for small $\Omega$ so that the levels are not significantly shifted. Figure 4 compares shifts of the one, three, and five photon processes for some typical parameters. Figure 4(a) shows these shifts for changes in $\Delta$ while figure 4(b) indicates the shifts for changes in $\Omega$. One can see that the shifts are much more
Figure 3 Various schemes allow the atom to be excited; (a) the usual one photon excitation where $\Delta = -kv$, (b) the three photon process, $\Delta = -3kv$, and (c) the five photon process, $\Delta = -5kv$. These transitions involve absorption of frequencies, $\omega^+_n$, and emission of frequencies, $\omega^-_n$.

\[ \omega^+_1 \quad \omega^+_3 \quad \omega^+_5 \]
\[ \omega^-_3 \quad \omega^-_5 \quad \omega^-_5 \]

one photon
three photon
five photon

($\Delta < 0$)
Figure 4  (a) Doppleron resonances occur at velocities which increase with $\Delta$ as shown. (b) Similarly, the resonant velocities increase with $\Omega$, the one photon process showing the largest shift. The velocities indicated below are the locations of the Doppleron peaks in figure 1 for various detunings and Rabi frequencies.

\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure4}
\caption{(a) $\Omega = 60\gamma$ \hspace{1cm} (b) $\Delta = 30\gamma$}
\end{figure}
strongly dependent on the detuning for the one photon process than for the higher order processes, which is apparent from equation 2.

The force due to the Dopplerons is similar to that of the RRP except that the net momentum exchanged can be more for the Dopplerons. In general, a momenta of \((2n+1)\hbar\kappa\) is imparted to the atom. The spontaneous emission rate will remain constant for these processes. However, the stimulated absorption rate for these events is finite and determined by \(\Omega\), such that the maximum force for higher order processes must fall off because of increasingly larger amounts of time necessary to excite the atom. Also, the probability of the multiphoton event goes as the product of the separate probabilities, therefore, larger order processes will have smaller transition probabilities. This effect is seen in figure 1(a), as \(\Omega\) is increased more Doppleron peaks in the force become visible and the relative sizes of these peaks changes.

The width of the Doppleron peaks decreases with \(n\) as can be understood by equation 2 and the preceding argument. The width of the one photon peak is power broadened, this is due to the effect of an increased stimulated emission rate for large \(\Omega\). The increased emission rate allows a larger spread of frequencies to be emitted. However, as the value of \(n\) increases this allowed spread is reduced by the factor of \(1/(2n+1)\), reducing the width of the force peak.

The resonances in the force can also be understood as a
result of Landau-Zener (LZ) like transitions between dressed states of the atom. In this picture, atoms make transitions between dressed levels non-adiabatically. As the atom moves through a field node, there is a finite probability it will change dressed states. This probability depends on the strength of the perturbation between the states and the velocity of the atom through the node. The LZ transitions lead to a coherence between the states which gives rise to oscillations in the dressed state populations. These oscillations are responsible for the sharp resonances in the force.
III. Apparatus

As shown above, the force on an atom in a standing wave can be divided into three basic features: the dipole force, the multiphoton processes and the remnants of the RRP. These effects are studied in two different experiments. The first experiment tests the collimation properties of the dipole force. The second experiment resolves the multiphoton resonances in the force. Both experiments are atomic beam experiments. A standing wave crosses an uncollimated atomic beam at a right angle. This geometry allows deflection of atoms in the beam using the velocity dependent force provided by the standing wave (SW). In this case the force is determined by the component of atomic velocity along the SW. Observation of the spatial collimation of the atomic beam requires measuring the transverse spatial distribution of the atoms. This is accomplished with a hotwire ionization detector. The multiphoton effects, however, are best understood by directly measuring the transverse velocity distribution. This can be accomplished by tuning a second probe laser across the resonance. The collimation effect will be emphasized for the particular choice of $\Omega$ and $\Delta$ that maximizes the dipole force peak. The Dopplerons will be most visible for a choice of $\Omega$ and $\Delta$ which emphasizes the Doppleron peaks. In general, the parameters for these two cases are
different. For either experiment the models that have been described involve a two-level system. In reality, no two-level atoms exist. However, one can approximate a two-level system by means of optical pumping. This chapter is divided into four main sections. The first section describes the atomic beam and vacuum system. The second section discusses the operation of the hotwire detector. The third section describes the lasers and optics used and their requirements. The final section describes the use of optical pumping to produce a two-level system for lithium.

III. A. Oven and Vacuum Chamber

An atomic beam of lithium can be produced by confining a sample of the metal in an oven. A small hole is put in the side of the oven to produce a beam. Provided the radius of the hole is on the order of or smaller than the mean free path of an atom in the vapor then the escaping atoms will provide an effusive beam. Effusive beams are to be contrasted with hydrodynamic flow in which the atoms collide in the hole and form a cloud around the exit. For many experiments these type of collisions are not desirable.

When the length of the hole, L, in the oven is small compared with its radius, r, the angular distribution of the flux from the hole will obey a cosine law.\[^{21}\]
\[ dQ = \frac{d\Omega}{4\pi} n\bar{v}A_s \cos \theta \]

with \( dQ \), the number of atoms per unit time through an aperture area, \( A_s \), into the solid angle \( d\Omega \) at \( \theta \) the angle from the normal to the hole. A detector of area \( A_d \) at a distance \( l \) from the aperture will see the number of atoms per unit time,

\[ I = \frac{A_d A_s}{4\pi l^2} n\bar{v} \cos \theta \]

As the length of the hole increases the flux becomes less and the angular distribution becomes narrower about \( \theta = 0 \). The deviation from the cosine law that we should expect for the case of our oven amounts to 10% at 10° from the normal. The significance of this will be discussed later when asymmetry in the data is discussed.

The oven was designed to be housed in a radiation shield which serves the dual purpose of insulating the rest of the chamber from the oven’s radiation and providing an efficient means of heating the oven. An aperture in the radiation shield allows part of the atomic beam into the vacuum chamber. This aperture effectively collimates the beam to within ± 9°. For lithium in our oven at a temperature of 500°C the most probable velocity is 1700 m/s. At 9° the corresponding transverse velocity is 300 m/s.
Our vacuum chamber operates well below $10^{-4}$ torr which is a much lower pressure than required to eliminate problems caused by interaction between the beam and the background gas. In general this condition is met when the mean free path of a beam atom in the background gas is larger than the chamber, corresponding to a pressure of $10^{-5}$ torr for our system.

III. B. Hotwire

Position dependent measurements of the atomic flux, downstream from the aperture, are made by means of a hotwire ionization detector. Neutral atoms adsorb onto a hot surface and will be ionized with some efficiency. The electron from the ionized atom is counted and the ion evaporates from the surface. This detector yields a current (ionized electrons) which is proportional to the incident number of atoms per unit time. The ratio, $\beta$, of ions to neutral atoms with ionization potential, $I$, evaporating from a hot surface (temperature $T$) with a work function, $\phi$, at thermal equilibrium is,

$$\beta = \frac{n_i}{n_n} = e^{-\frac{\delta}{kT}} \quad (3)$$

where $\delta = I - \phi$ and $n_i, (n_n)$ is the number of ions (neutrals) coming off the hotwire. In order to evaporate the neutrals and ions and increase the ionization fraction, the surface is
resisitively heated to a high temperature.

Considerations for detecting lithium demand a wire with a work function greater than or at least close to the first ionization potential of lithium, \( I = 5.392 \text{ eV} \). The best candidate is iridium which has, at best, an average work function of 5.4 eV. The reason for the average is because different crystal faces have different work functions, \( \Phi \). A typical wire of iridium will have a polycrystalline surface and so that an atom sees an average of \( \Phi \). Unfortunately, it is difficult to achieve a work function of 5.4 eV. More realistic iridium work functions fall below this value in a range of 5.0-5.4 eV. For detecting other atoms, the wire of choice is usually tungsten, however, its work function, \( \Phi_w = 4.6 \text{ eV} \), is 0.4-0.8 eV smaller than that of iridium. From equation 3 it is easy to see that such a difference in work function can cause an order of magnitude difference in detection efficiency.

Given our choice of Ir we are in the region of \( \delta > 0 \), therefore it appears to be advantageous to work at as high a temperature as possible. However, we found that increasing the temperature beyond a certain amount actually decreased the signal. We have no explanation for this effect.

Figure 5 shows the hotwire detection circuit. The voltage, \( V_{heat} \), provides the current necessary to heat the Ir wire. A battery was chosen for this source since it offers an extremely steady current and leakage of current to ground from
Figure 5 A voltage $V_{\text{heat}}$ is placed across a 0.005" Iridium wire heating the wire. Current generated by ionization of impinging atoms is collected by $V_{\text{bias}}$ and monitored by an electrometer.

\begin{align*}
\text{hotwire} \\
V_{\text{heat}} & \\
i_{\text{sig}} & \\
V_{\text{bias}} & \\
pA &
\end{align*}

$e^-$

$\rightarrow$

$\rightarrow$

$\rightarrow$
the source is eliminated. Current leakage is the ultimate source of error in measuring the current due to the atomic flux. A set of 9V batteries supply the bias voltage, $V_{\text{bias}}$. The excess electrons, generated as a result of the ionization of neutrals at the hotwire, produces a current in the bias circuit which is measured by an electrometer. The electrometer is capable of measuring currents smaller than 1 pA.

In order to achieve usable results with the Li-IR detection scheme is necessary to clean the hotwire once it is placed under vacuum. This is achieved by baking the wire at high temperature for a long time while a positive bias voltage is applied to repel ions. The net effect of the long term high temperature bakeout is to remove impurities from the surface of the wire. New wires require initial bakeout times of several days with a current through a five mil wire of 1.4 Amps. In addition, it is necessary to constantly bake the wires when the hotwire is not used. This technique is not foolproof. The steady-state signal varies with time, dropping off upon initial exposure to the beam with a rapid time constant of approximately 5 min, afterwards only slowly drifting. Measurements show the long term drift does not vary monotonically in time. For our experiment, it is sufficient that the drift over a half hour period be small compared to the signal size which is the case in general. A typical signal for an uncollimated beam and a heating current of 0.9 Amps is 500 pA, about what we expect for the estimated flux.
III. C. Lasers and Optics

The production of the force, in these experiments, requires an intense monochromatic standing-wave field. A stable, narrow bandwidth, tunable light source is also necessary as a velocity probe in the Doppleron experiment. The transition in lithium we want to excite is in the visible at 671 nm. These requirements can be met by a ring dye laser, which can produce linewidths below 1 MHz and single mode output powers of up to 1 W. The appropriate choice of dye for this case, DCM, offers a broad tunability over the required range of frequencies. Two dye lasers are used. The first, a Coherent 699-21, provides the standing wave (SW) while the second, a Spectra Physics 380D, provides the probe beam.

1. Collimation Beam

a. Frequency Stability

Frequency drifts of the SW during measurements will tend to smear out the data. The 699-21 has a long term drift of 50-100 MHz/hr (8-17 γ/hr). Changes in the detuning of this order cannot be tolerated, therefore, this laser is frequency locked to an external lithium vapor cell by means of saturated absorption spectroscopy. The use of acousto-optic modulators (AOM) to shift the frequency of the laser output allow the
laser frequency to be locked at some calibrated offset from the resonance frequency. By this method it is possible to lock the 699-21 to detunings in the range ±300 MHz (±50 γ) with a frequency deviation of less than γ. Another advantage of this system is that the detuning is read directly from the A/O modulation frequency.

b. Wavelength Measurement

Searching for the narrow resonance frequency (6 MHz) could be a very tedious process. Assistance in this measurement comes from a scanning Michelson interferometer wavemeter built after the design of Hall and Lee.23 The fringes from a known reference (in this case a stabilized He-Ne) are compared with the fringes from the unknown wavelength. The ratio of the number of fringes of the reference and unknown, for a given common scan length, yields the wavelength of the unknown. With this design an overall accuracy of one part in 10⁶ is obtained. This corresponds to an accuracy of ~500 MHz, which is sufficient to distinguish the 10 GHz separation of the 2P fine structure. Superior accuracy can be obtained by modifications to the system including interpolating between fringes. Relative accuracies exceeding 1:10⁶ have been achieved by Hall at JILA. An accuracy of 1:10⁶ is adequate for our purposes and provides reproducible locations of the resonances, allowing quick tuning to resonance.
c. Beam expansion

Extended interaction times between the atoms and the standing wave are achieved by expansion of the standing wave along the atomic beam. Usually cylindrical lenses are used to generate such a beam shape. Unfortunately, for large wavefronts, lenses introduce an unacceptable amount of distortion. An alternative is the use of a prism beam expander. A beam of light incident on a prism near Brewster's angle will expand upon exit from the other side for an appropriate choice of prism angle. For example, take the special case of a beam incident at Brewster's angle onto a Littrow prism. A Littrow prism is a right angle prism with Brewster's angle defining the rest of the triangle. In this case the magnification of the beam diameter is the ratio of the glass index of refraction to air, or the tangent of Brewster's angle. By adjusting the angle of incidence the magnification can be continuously varied. This type of beam expander is only efficient for beams polarized linearly in the plane of incidence, otherwise there is a significant power loss due to reflection. Fortunately, the 699-21 puts out a linearly polarized beam and a half-wave plate allows rotation of the axis of polarization to achieve beam expansions in any one direction.

d. Rabi Frequency Determination
The 699-21 is capable of up to one Watt continuous-wave output at 671 nm. Measurements of the power in the beam are made with a Coherent model 101 pyroelectric power meter, which has an absolute error of 1%. If the beam intensity profile is Gaussian then the waist can be determined by using a razor blade and micrometer. If \( \Delta x \) is the difference in positions of the razor blade when 0.3 and 0.7 of the total power is measured, then the beam waist, \( w \), is given by 1.91 \( \Delta x \). For an elliptically shaped beam, such as the expanded beam, the measurements along and perpendicular to the expansion can be made independently. In this case the intensity, \( I \), at the center of the Gaussian beam with total power, \( P_o \), is,

\[
I = \frac{2P_o}{\pi w_h w_v}
\]

where \( w_v \) (\( w_h \)) is the vertical (horizontal) beam waist. Finally, the Rabi frequency, \( \Omega \), is given as,

\[
\Omega = \frac{dE_o}{\hbar} = \gamma \sqrt{\frac{I}{I_{sat}}} = \gamma \sqrt{\frac{2P_o}{\pi w_h w_v I_{sat}}}
\]

for a saturation intensity, \( I_{sat} \), a travelling wave of field strength, \( E_o \), and induced dipole moment, \( d \). The saturation intensity is the intensity of light for which \( \Omega = \gamma \). The Rabi frequency in a standing wave is larger by a factor of two.
2. Probe Beam

The probe beam measures the transverse velocity distribution of the atomic beam. Atoms in resonance with the probe beam will fluoresce isotropically due to spontaneous emission which can be detected. An atom moving with a velocity, $\mathbf{v}$, with respect to the photon wave vector, $\mathbf{k}$, will see the Doppler shifted frequency,

$$\omega = \omega_L - k \cdot \mathbf{v}$$

Given a narrow linewidth laser, only those atoms with velocities such that $\omega \approx \omega_0$, where $\omega_0$ is the resonance frequency, will scatter light and contribute to the fluorescence. Therefore, the fluorescence intensity for a given laser frequency, is a relative measure of the number of atoms with a velocity component along the probe beam satisfying the above resonance condition. For a typical atomic beam, the transverse velocity distribution is Gaussian and so a plot of fluorescence as a function of laser detuning should be approximately Gaussian. If the velocity distribution has been altered, for example, by the force produced as the atoms traverse the standing wave, then the velocity distribution will be changed. This change will become manifest in the fluorescence signal. This technique is limited by measurements of the probe frequency, the probe linewidth and the natural linewidth of the transition excited.
a. Frequency Stability

The frequency stability of the probe is somewhat less demanding. The long term stability is not as important as the short term jitter and linearity of scan. The probe laser, a Spectra Physics 380D, has a short term jitter of less than 1 MHz. The linearity of scan is specified at 1% over the maximum scan range of 30 GHz. With these specifications, the limit on resolution comes from the natural linewidth of the transition which is \((2\pi)5.83\) MHz.\(^2\) This should allow the resolution of velocity structure greater than 4 m/s \((1 \gamma/k)\). This resolution must be taken into consideration when choosing the SW parameters. The force plots of figure 1 indicate that with a resolution of \(\gamma/k\) it will be difficult to resolve higher order Doppleron resonances.
b. Scan Width Calibration

In order to measure the relative positions and widths of structure, the probe frequency scan width must be calibrated. This is accomplished with a scanning Fabry-Perot interferometer (FPI). The transmission signal of a scanning FPI is used to observe the mode structure of the probe. The signal is detected by a photodiode and gives a trace on an oscilloscope. Figure 6 shows a typical set of FPI fringes for a single mode dye laser. The fringe spacing for the two FPI modes visible at a given input frequency is the free spectral range of the FPI, in this case ~1600 MHz. By proper adjustment of the scope, the horizontal divisions can be calibrated in frequency. By observing the difference in positions of the transmission peaks at the beginning and end of a probe scan, it is possible to determine the absolute frequency width of the scan. This technique is limited by thermal drifts which can be large for long scan times. At the end of the scan the laser is quickly returned to the starting frequency. The drift of the transmission peaks during this rewind time is negligible and therefore it becomes advantageous to compare the peak locations between the end of the scan and the end of the rewind.
Figure 6 A typical set of FPI fringes at the beginning and end of a scan. The positions can be calibrated by the known free spectral range of the FPI (~1600MHz), allowing measurements of the scan range.

\[
\text{scan range} = \frac{a-b}{b} \times \text{FSR}
\]

FPI transmission

\[
\text{FSR}
\]

start of scan

end of scan

horizontal axis
(relative cavity spacing)
III. D. Optical Pumping of Lithium

Optical pumping describes a process of generating a non-thermal distribution among the various energy levels of an atom. In this case the pumping is achieved by means of irradiation by a laser of appropriate polarization. The relevant energy levels of lithium are shown in figure 7. If a directional right circularly polarized beam is tuned to the $2P_{3/2}, F=3 \leftrightarrow 2S_{1/2}, F=2$ resonance it will induce a quantization axis onto the atoms. The selection rules for dipole absorption are $\Delta F = 0, \pm 1, \Delta m_F = \pm 1$. However, the selection rules for spontaneous emission are the same for $F$ but for $m_F$ become $\Delta m_F = 0, \pm 1$. If the laser is sufficiently narrow, atoms originally in the lower ground state ($2S_{1/2}, F=1$) will remain there. Atoms in the upper ground state will make transition between $2S_{1/2}$ and $2P_{3/2}$ as shown in figure 7. After some amount of time atoms will either be pumped into the cycling transition ($2P_{3/2}, F=3, m_F=3 \leftrightarrow 2S_{1/2}, F=2, m_F=2$) or fall into the lower ground state. The cycling transition constitutes our two-level system. The ground state hyperfine structure is split by $v_{hf}=803$ Mhz. However, for an intense laser the energy levels are broadened and some of the light will excite atoms from the lower ground state. This will increase the steady state population in the cycling transition. A second beam can also be used to excite the lower ground state atoms with similar results.
Figure 7 Relevant lithium energy levels for generation of a two-level system in our experiment. The solid lines indicate possible stimulated processes for right hand circularly polarized light and the dotted lines indicate spontaneous transitions.
IV. Experiments

This chapter describes the experiments which probe the features of the force on an atom in a strong standing wave. Two experiments are described. The first experiment tests the dipole force by measuring the collimation of an atomic lithium beam after traversing an orthogonal standing wave. The data is shown to agree with expected results obtained by integrating the equation of motion of the atom in the standing wave for a relevant range of transverse velocities. The second experiment resolves the structure in the force due to multiphoton (Doppleron) processes. This structure is measured by observing the transverse velocity distribution of the atomic beam after the interaction region. Structure corresponding to Doppleron processes of up to nine-photon is resolved. These data are in good agreement with fit calculations obtained similar to those above.

IV. A. Laser Collimation of an Atomic Beam

In this experiment, atoms with a given transverse velocity have their transverse velocities reduced, effectively collimating the atomic beam. This arrangement is shown in figure 8(a). The force plots of figure 1 show that for a given value of $\Omega$, positive detunings cools atoms with
Figure 8 Experimental setups to obtain (a) collimation and (b) Doppleron data. Light from a dye laser is expanded along the z-axis for increased interaction times, circularly polarized to optically pump into a two-level system, and reflected back on itself to produce a standing-wave. (a) The effects on the spatial distribution are measured by scanning a hotwire along the x-axis 30 cm downstream from the oven. (b) The effects on the velocity distribution are detected by scanning a probe laser across the atomic resonance and detecting the fluorescence.
velocities less than some critical velocity, $v_0$, determined by the zero crossing of the force. Atoms with $v > v_0$ will be heated. For $\Delta > 0$, atoms with $v_x < v_0$ will accumulate about $v_x = 0$ while atoms with $v_x > v_0$ will be pushed out to higher $v_x$. When the sign of $\Delta$ is reversed, the atoms near $v_x = 0$ are heated to the zero crossing at $v_x = v_0$, where they accumulate. Likewise, atoms with $v_x > v_0$ are cooled and accumulate at $v_x = v_0$.

1. Collimation Data

The amount of collimation depends on the time the atom spends in the standing wave. Longer times are provided by expanding the SW along the atomic flux using the prism beam expander. The spatial distribution of the flux is detected downstream using a hotwire ionization detector. The standing wave is generated by reflecting the collimation beam back on itself with a mirror. Losses in intensity through the anti-reflection coated vacuum windows are less than 1% which ensures that $\Omega$ remains the same in or out of the chamber. Circular polarization for optical pumping is achieved by using a quarter wave plate.

The experiment was performed using $\Omega = 50 \gamma$ with an interaction length of $\sim 10$ mm. The collimation signal was maximized by tuning the SW frequency across the atomic $2S_{1/2} - 2P_{3/2}$ resonance while the hotwire signal at 30 cm from the oven
is monitored. The hotwire was then scanned over 25 mm with the SW at the frequency yielding the largest change from the uncollimated signal. Several other hotwire scans are also performed at various SW frequencies. All these scans show structure consistent with the above discussions, however, the signal remained small.

At this point a secondary laser beam was added which served to re-pump atoms from the previously unreachable $2S_{1/2}$, $F=1$ ground state into states accessible to the SW. This secondary beam, tuned 803.5 Mhz (the ground state hyperfine splitting) above the SW, was generated using an acousto-optic (A/O) modulator to frequency shift a small fraction of the SW power split off from laser 1. The net effect of this second beam is to increase the overall fraction of atoms which can contribute to the collimation signal. Data obtained in this manner show greater than 100% increase over the uncollimated flux. Figure 9 shows the data obtained for various detunings. The actual detunings corresponding to the different data are not known, however, the relative detunings and the absolute sign are known. Figure 9(a) shows hotwire scans for various negative detunings. The curves show two peaks corresponding to the location of the critical velocities. For smaller $|\Delta|$ the two peaks come together since $v_0$ decreases with $|\Delta|$. Finally, the two peaks can no longer be resolved. The plots also show areas where the flux has been depleted due to heating and cooling atoms out of those regions. Figure 9(b)
Figure 9: (a) shows data obtained by the hotwire for various \( \Delta < 0 \). Atoms collect about the critical velocities, \( v_c \). (b) shows the case for \( \Delta > 0 \) where atoms collect at \( v_x = 0 \) and are depleted from the \( v_x = v_c \) regions. The narrowest peak in (b) corresponds to a collimation of \( \pm 1.6 \) mrad at full-width half-max of the central peak.
is a similar set of plots for positive detunings. In these plots there is only one central peak, corresponding to atoms with $v_n < v_0$ being cooled by the dipole force. The depletion regions due to heating are the cumulative effects of the different order Dopplerons.

2. Model of the Collimation Data

This data was modeled by a program which calculated the expected atomic intensity at the hotwire for a given distribution of atoms traversing a region of standing wave using forces calculated as described above. The calculation propagates the initial velocity distribution through the region of the standing wave. The resulting velocity and position distribution is propagated to the hotwire position where the positions are histogrammed into bins corresponding to the resolution of the hotwire. The bins are then averaged over the aperture diameter producing the final transverse intensity distribution. The expected distribution was calculated for $\Omega = 50 \, \gamma$ and various detunings. The results are shown in figure 10 and show good qualitative agreement with the general features seen in the experimental data. The broadening of the features in the experimental data is most likely due to the atoms sampling a large range of Rabi frequencies as they pass through the Gaussian beam waist of the SW. The difference in the magnitude of the dips and peaks
Figure 10 Calculated spatial distributions of the flux for $\Delta>0$ (a) and $\Delta<0$ (b). The fit has narrower features than the data due to the random nature of the force and frequency jitter in the collimation laser. The shape of the data qualitatively agrees with the fit. Atoms accumulate about $v_x=v_o$ for $\Delta<0$ and about $v_x=0$ for $\Delta>0$. 

(a) $\Omega = 50\gamma$

(b) $\Omega = 50\gamma$
in the intensity distribution of the data and the fit is most likely due to the random nature of the force and frequency jitter of the standing wave. In the calculations atoms will remain at stable zero crossings. In reality, the atoms see fluctuations in the force. For example, this would tend to reduce and broaden the peak in the flux at $v_x=0$ for blue detunings.

The best collimation is achieved using blue detunings since the flux peaks about $v_x=0$. The data of figure 9(b), for $\Delta > 0$, indicate a full collimation angle at FWHM of better than 3.0 mrad (0.2 degrees) with an increase in flux of 100% over the uncollimated beam.

Finally, it can be seen by the data in figures 9 and 10 that the Doppleron resonances are not resolved. If the atomic beam possessed one longitudinal velocity this would not be the case. However, the wide distribution of the longitudinal velocities exiting from the oven smears out the spatial locations of the velocity dependent Dopplerons. A given Doppleron exists at a particular transverse velocity. The location of this Doppleron spatially along the x-axis will depend on both the transverse and longitudinal velocities of the particular atom. The large spread in longitudinal velocities acts to continuously overlap the Doppleron information in space. Clearly, a detection mechanism which is sensitive to velocities would be better suited for resolving Dopplerons. This is the case of Doppler spectroscopy and is
the subject of the next section. Alternatively, the atomic flux could be filtered using a velocity selector which would allow only a narrow width of longitudinal velocities through, thus permitting the observation of the Doppleron resonances with the spatial detection scheme.

IV. B. Observation of Doppleron Resonances

This experiment investigates the redistribution of transverse atomic velocities by a standing wave, specifically, the effect of the Doppleron resonances is probed. From figure 1 the Doppleron resonances provide heating for $A > 0$ and cooling for $A < 0$. Atoms with $v_x$ coinciding with a Doppleron resonance will be quickly heated or cooled to a velocity with a smaller force. For a very long interaction time the atoms will either heat out to large velocities or accumulate at the stable zero crossings of the force. For smaller times the atoms will be repulsed from the regions in $v_x$ where there are Doppleron resonances. This redistribution is detected by means of Doppler spectroscopy as described above.

1. Doppleron Data

In this experiment the standing wave and atomic beam are setup as in the collimation experiment (without the A/0 modulated beam). A second weak beam parallel to the standing
wave serves as the velocity probe (figure 8(b)). The standing wave intensity of 4.7 W/cm² yields Ω=60γ. For this Ω a detuning of ±30γ is chosen to emphasize the Dopplerons. The lowest order Dopplerons are separated by more than a linewidth for this detuning. This is important since the resolution of the velocity probe is limited by γ, which translates to 4 m/s in velocity. The SW is clipped on both sides to minimize the variation of Ω to less than 10% of the maximum value. The shifts of the Dopplerons due to the variation of Ω introduces an uncertainty of ~5% in the location of the one photon peak. The collimation laser is locked to a lithium vapor cell to eliminate drifts which will smear out the positions of the Dopplerons. The uncertainty of the lock frequency of the SW is less than 1.5 γ. This introduces an error in the Doppleron resonance locations of 1.5 γ for the one photon process, 0.5 γ for the three, and 1.5γ/l, for the 1st order Doppleron. The probe is scanned across the 2S₁/₂-2P₃/₂ transition while the fluorescence is collected by a photodiode. The width of the scan is determined using the FPI as described above. The uncertainty in this width is about 3%.

Figure 11 shows the data obtained for the above conditions. The solid lines are experimental values and the dotted lines are fits. The arrows and numbers indicate the locations of the dips in the velocity distribution due to the Doppleron resonance peaks. Figure 11(a) shows the data for positive detuning. The large central peak is due to the atoms
Figure 11 Data obtained by spectroscopy of collimated atoms. Solid line are data points and dotted lines are fits. (a) shows data for \( \Delta > 0 \). Atoms are heated by the Doppleron resonances to regions of smaller force resulting in dips in the distribution at these velocities. (b) shows the case for \( \Delta < 0 \). The fits are in good agreement with the data including the modelled asymmetry.
collimated by the dipole force. Atoms are heated by the Doppleron peaks and bunch up at velocities between the peaks. Figure 11(b) shows the data for negative detuning. The central dip results from the dipole force heating atoms out of the $v_z=0$ region. The atoms are cooled out of the region of the peaks.

The asymmetry in the data is due to slight misalignment from $90^\circ$ for the angle between the atomic beam axis and the SW. The angular distribution of our atomic beam is sharp due to the oven radiation shield aperture. When the beam axis and the SW are not at a right angle the angular distribution favors one side of the Doppler profile, causing the asymmetry.

2. Fit of the Doppleron Data

The fit calculations are similar to those described above for the collimation experiment, except the transverse velocity is histogrammed into velocity bins after the standing wave. This fit uses the initial transverse velocity distribution (obtained without the SW) as an input. The bins are then averaged over the Lorentzian linewidth of the probe providing the final velocity distribution. The overall normalization and zero velocity are the only input fitting parameters. The frequency scale of the experimental data was scaled within the experimental error to improve the agreement with the theory. The fit shows good agreement with the experimental
distribution including the relative locations of the dips in the distribution due to the Doppleron peaks. The central peak (dip) for the case of $\Delta > 0$ ($\Delta < 0$) is lower (higher) for the experimental data than for the fit due to the fluctuations in the dipole force, and the fluctuations of the lasers. The fit disagrees with the experimental data in the region of the one photon peaks. This disagreement is likely due to atoms traversing a weaker region of the standing wave radiation, which would give narrower one-photon resonances. Finally, the slight asymmetry in the data is modeled by adjusting the relative velocity zero of the input distribution, showing good agreement with the data.
V. Conclusion

The force on an atom moving in a standing wave can be large enough to collimate a thermal atomic beam. This means the collimation increases the intensity in the central region and provides a narrow divergence angle for this central beam. Experiments show structure in the spatial distribution of atoms collimated in this fashion. Comparison to calculations show that this structure is due to the dipole force an atom feels when moving in a standing wave. The data and calculations indicate that this force can be much larger than that due to RRP. Other experiments show the presence of structure in the velocity distribution of these collimated atoms. This structure is shown to be in agreement with that predicted by theory, indicating that it is produced by multiphoton (Doppler-on) resonances in the force. The strength of the dipole force will allow atoms in a beam to be slowed longitudinally over distances much shorter than achieved in previous beam slowing techniques employing RRP. These techniques should prove useful for generating intense sources of cold atoms for trapped atom experiments.
VII. References


