

# Ph 76 ADVANCED PHYSICS LABORATORY

## – ATOMIC AND OPTICAL PHYSICS –

### Saturated Absorption Spectroscopy

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#### I. BACKGROUND

One of the most important scientific applications of lasers is in the area of precision atomic and molecular spectroscopy. Spectroscopy is used not only to better understand the structure of atoms and molecules, but also to define standards in metrology. For example, the second is defined from atomic clocks using the 9192631770 Hz (exact, by definition) hyperfine transition frequency in atomic cesium, and the meter is (indirectly) defined from the wavelength of lasers locked to atomic reference lines. Furthermore, precision spectroscopy of atomic hydrogen and positronium is currently being pursued as a means of more accurately testing quantum electrodynamics (QED), which so far is in agreement with fundamental measurements to a high level of precision (theory and experiment agree to better than a part in  $10^8$ ). An excellent article describing precision spectroscopy of atomic hydrogen, the simplest atom, is attached (Hänsch *et al.* 1979). Although it is a bit old, the article contains many ideas and techniques in precision spectroscopy that continue to be used and refined to this day.

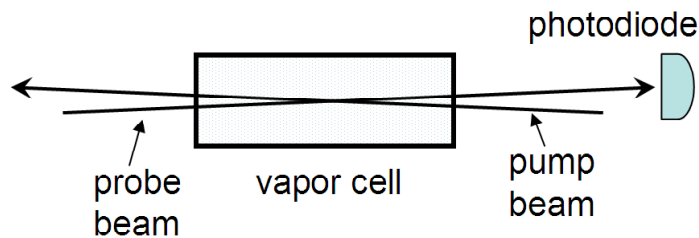


Figure 1. The basic saturated absorption spectroscopy set-up.

**Qualitative Picture of Saturated Absorption Spectroscopy – 2-Level Atoms.** Saturated absorption spectroscopy is one simple and frequently-used technique for measuring narrow-line atomic spectral features, limited only by the natural linewidth  $\Gamma$  of the transition (for the rubidium D lines  $\Gamma \approx 6$  MHz), from an atomic vapor with large Doppler broadening of  $\Delta\nu_{Dopp} \sim 1$  GHz. To see how saturated absorption spectroscopy works, consider the experimental set-up shown in Figure 1. Two lasers are sent through an atomic vapor cell from opposite directions; one, the “probe” beam, is very weak, while the other, the “pump” beam, is strong. Both beams are derived from the same laser, and therefore have the same frequency. As the laser frequency is scanned, the probe beam intensity is measured by a photodetector.

If one had 2-level atoms in the vapor cell, one might record spectra like those shown in Figure 2. The upper plot gives the probe beam absorption without the pump beam. Here one sees simple Doppler-

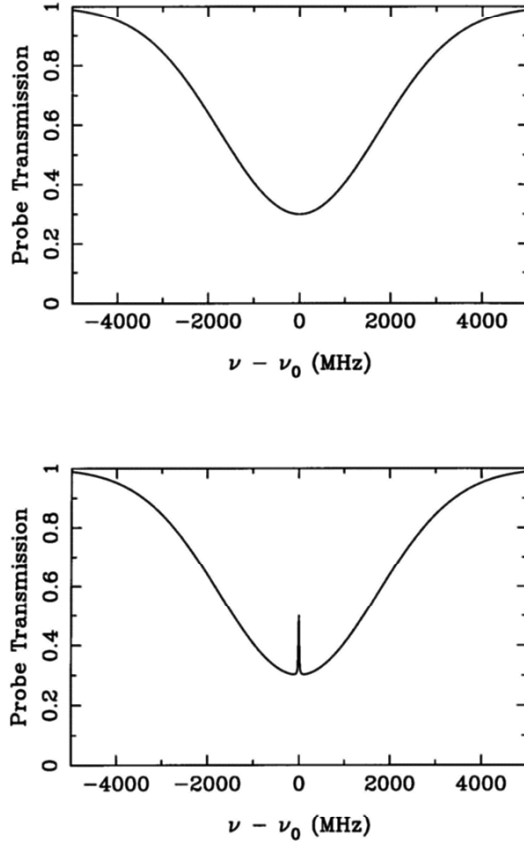


Figure 2. Probe absorption spectra for 2-level atoms, both without (upper) and with (lower) the pump beam.

broadened absorption; in our case the Doppler width is much larger than the natural linewidth,  $\Delta\nu_{Dopp} \gg \Gamma$ , and the optical depth of the vapor is fairly small  $\tau(\nu) \lesssim 1$  (the transmitted fraction of the probe is  $e^{-\tau(\nu)}$ , which defines the optical depth;  $\tau$  is proportional to the atomic vapor density and the path length), so the probe spectrum is essentially a simple Gaussian profile.

The lower plot in Figure 2 shows the spectrum with the pump beam, showing an additional spike right at the atomic resonance frequency. The reason this spike appears is as follows: If the laser frequency is  $\nu_0 - \Delta\nu$ , then the probe beam is absorbed only by atoms moving with longitudinal velocity  $v \approx c\Delta\nu/\nu_0$ , moving toward the probe beam. These atoms see the probe beam blueshifted into resonance; other atoms are not in resonance with the probe beam, and so they do not contribute to the probe absorption. These same atoms see the pump beam red-shifted further from resonance (since the pump beam is in the opposite direction) so they are unaffected by the pump beam. Thus for laser frequencies  $\nu \neq \nu_0$ , the probe absorption is the same with or without the pump beam. However if  $\nu = \nu_0$ , then atoms with  $v = 0$  contribute to the probe absorption. These  $v = 0$  atoms also see an on-resonance pump beam, which is strong enough to keep a significant fraction of the atoms in the excited state, where they do not absorb the probe beam (in fact they increase the probe beam intensity via stimulated emission). Thus at  $\nu = \nu_0$  the probe absorption is *less* than it was without the pump beam. (If the pump beam had infinite intensity, half of the atoms would

be in the excited state at any given time, and there would be identically zero probe absorption. One would say these atoms were completely “saturated” by the pump beam, hence the name saturated absorption spectroscopy.) The advantage of this form of spectroscopy should be obvious . . . one can measure sharp Doppler-free features in a Doppler-broadened vapor.

**Qualitative Picture of Saturated Absorption Spectroscopy – Multi-level Atoms.** If the atoms in the absorption cell had a single ground state and two excited states (typically an electronic level split by the hyperfine interaction), and the separation of the excited states was less than the Doppler width, then one would see a spectrum like that shown in Figure 3. The peaks on the left and right are ordinary saturated absorption peaks at  $\nu_1$  and  $\nu_2$ , the two resonance frequencies. The middle peak at  $(\nu_1 + \nu_2)/2$  is called a “cross-over resonance.” If you think about it for a while you can see where the extra peak comes from. It arises from atoms moving at velocities such that the pump is in resonance with one transition, and the probe is in resonance with the other transition. If you think about it a bit more you will see there are *two* velocity classes of atoms for which this is true – atoms moving toward the pump laser, and away from it.

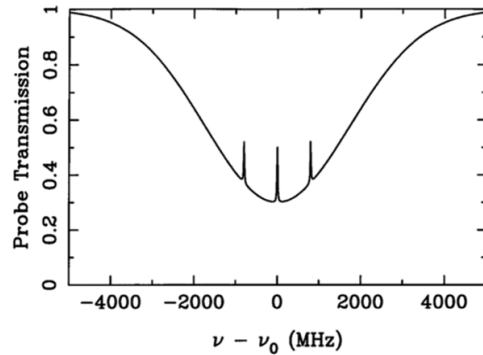


Figure 3. Saturated absorption spectrum for atoms with a single ground state and two closely spaced excited states.

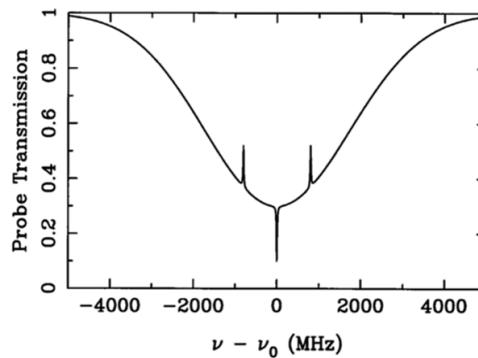


Figure 4. Saturated absorption spectrum for atoms with a single excited state that can decay into either of two closely spaced ground states.

If the atoms in the vapor cell had a single excited state but two hyperfine ground states (we call them both “ground” states because neither can decay via an allowed transition), and the separation of the ground states was less than the Doppler width, then one might see a spectrum like in Figure 4. The extra cross-over dip results from a phenomenon called “optical pumping,” which occurs because atoms in the excited state can decay into either of the two stable ground states. Thus if atoms are initially in ground state  $g1$ , and one shines in a laser that excites  $g1 \rightarrow e$ , atoms will get excited from  $g1 \rightarrow e$  over and over again until they once spontaneously decay to  $g2$ , where they will stay. The state  $g2$  is called a “dark state” in this case, because atoms in  $g2$  are not affected by the laser. We see that a laser exciting  $g1 \rightarrow e$  will eventually optically pump all the atoms into  $g2$ .

To see how optical pumping produces the extra crossover dip, remember that only the pump laser can optically pump – the probe laser is by definition too weak. Also remember the atoms in the cell are not in steady state. When they hit the walls they bounce off about equally distributed in both ground states, and the optical pumping only operates for a short period of time as the atoms travel through the laser beams. If you think about it a while you can see there are two velocity classes of atoms that are responsible for the dip. For one velocity class the pump laser excites  $g1 \rightarrow e$ , which tends to pump atoms into  $g2$ . Then the probe laser, which excites  $g2 \rightarrow e$  for these atoms, sees extra absorption. For the other velocity class the pump laser excites  $g2 \rightarrow e$ ,  $g1$  gets overpopulated, and again the probe laser (which now excites  $g1 \rightarrow e$  for these atoms) sees more absorption.

**Quantitative Picture of Saturated Absorption Spectroscopy – 2-Level Atoms.** One can fairly easily write down the basic ideas needed to calculate a crude saturated absorption spectrum for 2-level atoms, which demonstrates much of the underlying physics. The main features are: 1) the transmission of the probe laser beam through the cell is  $e^{-\tau(\nu)}$ ,  $\tau(\nu)$  is the optical depth of the vapor; 2) the contribution to  $\tau(\nu)$  from one velocity class of atoms is given by

$$d\tau(\nu, v) \sim (P_1 - P_2)F(\nu, v)dn(v)$$

where  $P_1$  is the relative population of the ground state,  $P_2$  is the relative population of the excited state ( $P_1 + P_2 = 1$ ),

$$dn \sim e^{-mv^2/2kT} dv$$

is the Boltzmann distribution (for  $v$  along the beam axis), and

$$F(\nu, v) = \frac{\Gamma/2\pi}{(\nu - \nu_0 + \nu_0 v/c)^2 + \Gamma^2/4}$$

is the normalized Lorentzian absorption profile of an atom with natural linewidth  $\Gamma$ , including the Doppler shift. Putting this together, we have the differential contribution to the optical depth, for laser frequency  $\nu$  and atomic velocity  $v$ :

$$d\tau(\nu, v) = \tau_0 \frac{\nu_0}{c} (P_1 - P_2) F(\nu, v) e^{-mv^2/2kT} dv.$$

The overall normalization comes in with the  $\tau_0$  factor, which is the optical depth at the center of resonance line, *i. e.*  $\tau_0 = \int d\tau(\nu_0, v)$  with no pump laser (the integral is over all velocity classes). 3) The populations of the excited and ground states are given by  $P_1 - P_2 = 1 - 2P_2$ , and

$$P_2 = \frac{s/2}{1 + s + 4\delta^2/\Gamma^2}$$

where  $s = I/I_{sat}$  and  $\delta = \nu - \nu_0 - \nu_0 v/c$ .  $I_{sat}$  is called the saturation intensity (for obvious reasons ... if you consider the above formula for  $P_2$  with  $\delta = 0$ ,  $P_2$  “saturates”  $P_2 \rightarrow 1/2$  as  $I/I_{sat} \rightarrow \infty$ ). The value of

$I_{sat}$  is given by

$$I_{sat} = 2\pi^2 hc\Gamma/3\lambda^3.$$

For the case of rubidium,  $\Gamma \approx 6$  MHz, giving  $I_{sat} \approx 2$  mW/cm<sup>2</sup>.

The underlying physics in points (1) and (2) should be recognizable to you. Point (3) results from the competition between spontaneous and stimulated emission. To see roughly how this comes about, write the population rate equations as

$$\begin{aligned}\dot{P}_1 &= \Gamma P_2 - \alpha I(P_1 - P_2) \\ \dot{P}_2 &= -\Gamma P_2 + \alpha I(P_1 - P_2)\end{aligned}$$

where the first term is from spontaneous emission, with  $\Gamma$  equal to the excited state lifetime, and the second term is from stimulated emission, with  $\alpha$  a normalization constant. Note that the stimulated emission is proportional to the intensity  $I$ . In the steady-state  $\dot{P}_1 = \dot{P}_2 = 0$ , giving

$$P_2 = \frac{\alpha I/\Gamma}{1 + 2\alpha I/\Gamma}$$

The term  $\alpha I/\Gamma$  corresponds to the  $s/2$  term above (note  $I_{sat}$  is proportional to  $\Gamma$ ). A more complete derivation of the result, with all the normalization constants, is given in Milonni and Eberly (1988), and in Cohen-Tannoudji *et al.* (1992), but this gives you the basic idea.

Assuming a fixed vapor temperature, atomic mass, etc., the saturated absorption spectrum is determined by two adjustable external parameters, the pump intensity  $I_{pump}$  and the on-resonance optical depth  $\tau_0$ . The latter is proportional to the vapor density inside the cell. Figure 5 shows calculated spectra at fixed laser intensity for different optical depths, and Figure 6 shows spectra at fixed optical depth for different laser intensities.

In Figure 5 one sees mainly what happens when the vapor density is increased in the cell. At low densities the probe absorption is slight, with a Gaussian profile, and the absorption increases as the vapor density increases. At very high vapor densities the absorption profile gets deeper and broader. It gets broader simply because the absorption is so high near resonance that the probe is almost completely absorbed; for greater vapor densities the probe gets nearly completely absorbed even at frequencies fairly far from resonance; thus the width of the absorption profile appears broader. The saturated-absorption feature in Figure 5 does pretty much what you would expect. The probe absorption is reduced on resonance, due to the action of the pump laser. At very high vapor densities the saturated-absorption feature becomes smaller. This is because while the pump laser reduces the absorption, it doesn't eliminate it; thus at high vapor densities the probe is nearly completely absorbed even with the pump laser. The moral of this story is that the vapor density shouldn't be too low or high if you want to see some saturated-absorption features.

In Figure 6 one sees that if the pump intensity is low, the saturated-absorption feature is small, as one would expect. For larger pump intensities the feature grows in height and width. The width increases because at high laser intensities the effect of the pump laser saturates on resonance, and continues to grow off resonance; thus the width of the feature increases, an effect known as "power broadening."

Finally, it should be noted that calculating the saturated absorption spectrum for real atoms, which must include optical pumping, many different atomic levels, atomic motion in the vapor cell, and the polarization of the laser beams, is considerably more subtle. A recent paper by Schmidt *et al.* (1994) shows much detailed data and calculations for the case of cesium.

**Problem 1.** Show that  $\tau_0 = \int d\tau(\nu_0, \nu)$  when the pump laser intensity is zero, from the formula above.

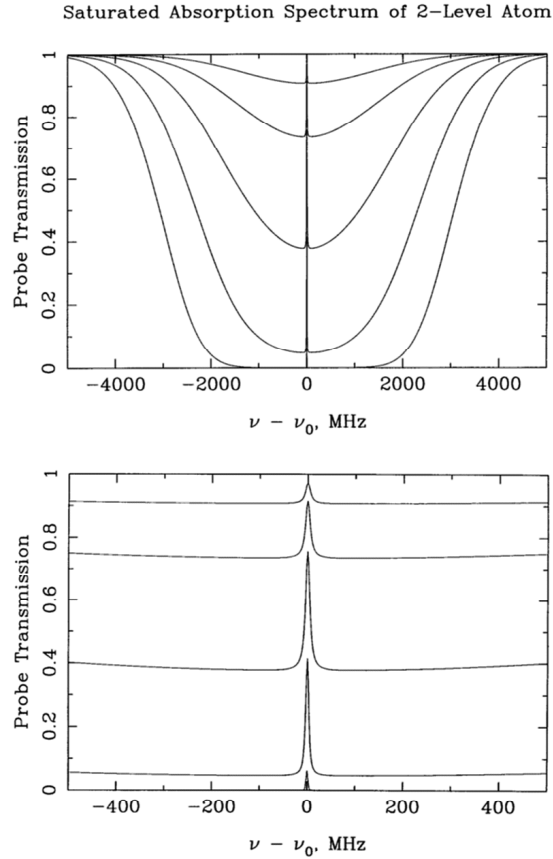


Figure 5. Calculated saturated-absorption spectra for two-level atoms, for  $(\tau, I/I_{sat}) = (0.1, 10), (0.316, 10), (1, 10), (3.16, 10),$  and  $(10, 10)$ . The two plots show the same spectra with the frequency axis at different scales. Note the overall Doppler-broadened absorption, with the small saturated-absorption feature at line center.

Hint: the integral is simplified by noting that  $\Gamma \ll \Delta\nu_{Dopp}$ .

**Problem 2.** The above calculations all assume that the pump laser has the same intensity from one end of the cell to another. This is okay for a first approximation, but calculating what really happens is an interesting problem. Consider a simple laser beam (the pump) shining through a vapor cell. If the laser intensity is weak, and the atoms are all pretty much in the ground state, then the laser intensity changes according to the equation  $dI/dx = -\alpha I$ , where  $\alpha = \alpha(\nu)$  depends on the laser frequency, but not on position inside the cell ( $\alpha^{-1}$  is called the absorption length in this case). This equation has the solution  $I(x) = I_{init}e^{-\alpha x}$ , where  $I_{init}$  is the initial laser intensity. The transmission through the cell,  $e^{-\alpha L}$ , where  $L$  is the length of the cell, is what we called  $e^{-\tau}$  above.

Your job in this problem is to work out what happens when the input laser beam is not weak, and thus we cannot assume that the atoms are all in the ground state. In this case  $\alpha = \alpha(\nu, x)$ , which makes the differential equation somewhat more interesting. Assume the laser is on resonance for simplicity. Then the attenuation coefficient at any position  $x$  is proportional to  $P_1 - P_2$ , which in turn is proportional to

### Saturated Absorption Spectrum of 2-Level Atom

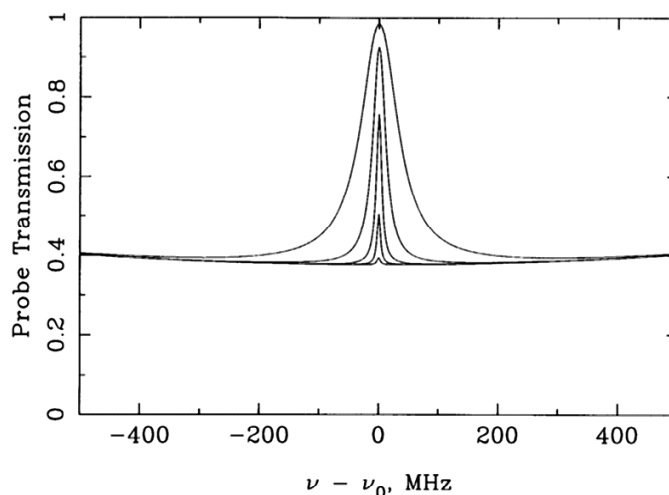


Figure 6. Calculated saturated-absorption spectra for two-level atoms, for  $(\tau, I/I_{sat}) = (1, 0.1), (1, 1), (1, 10), (1, 100),$  and  $(1, 1000)$ . Note at large laser intensities the saturated absorption feature is “power broadened” as the line saturates.

$1/(1 + s)$ . Thus we have  $\alpha(\nu_0, x) = \alpha_0/(1 + s(x))$ . In the weak beam limit  $I \ll I_{sat}$  this reduces to our previous expression, so  $\alpha_0 = \tau_0/L$ . Write down an expression which relates the saturation parameter of the laser as it exits the cell  $s_{final}$ , the saturation parameter at the cell entrance  $s_{initial}$ , and the weak-limit optical depth  $\tau_0$ . Check your expression by noting in the limit of finite  $\tau_0$  and small  $s$  you get  $s_{final} = s_{initial}e^{-\tau_0}$ . If  $\tau_0 = 100$ , how large must  $s_{initial}$  be in order to have a transmission of  $1/2$  (*i. e.*  $s_{final} = s_{initial}/2$ )?

**Atomic Structure of Rubidium.** The ground-state electronic configuration of rubidium consists of closed shells plus a single 5s valence electron. This gives a spectrum which is similar to hydrogen (see attached Scientific American article). For the first excited state the 5s electron is moved up to 5p. Rubidium has two stable isotopes:  $^{85}\text{Rb}$  (72 percent abundance), with nuclear spin quantum number  $I = 5/2$ , and  $^{87}\text{Rb}$  (28 percent abundance), with  $I = 3/2$ .

The different energy levels are labeled by “term states”, with the notation  $^{2S+1}L'_J$ , where  $S$  is the spin quantum number,  $L'$  is the spectroscopic notation for the angular momentum quantum number (*i. e.* S, P, D, ..., for orbital angular momentum quantum number  $L = 0, 1, 2, \dots$ ), and  $J = L + S$  is the total angular momentum quantum number. For the ground state of rubidium  $S = 1/2$  (since only a single electron contributes), and  $L = 0$ , giving  $J = 1/2$  and the ground state  $^2S_{1/2}$ . For the first excited state we have  $S = 1/2$ , and  $L = 1$ , giving  $J = 1/2$  or  $J = 3/2$ , so there are two excited states  $^2P_{1/2}$  and  $^2P_{3/2}$ . Spin-orbit coupling lifts splits the otherwise degenerate  $P_{1/2}$  and  $P_{3/2}$  levels. (See any good quantum mechanics or atomic physics text for a discussion of spin-orbit coupling.)

The dominant term in the interaction between the nuclear spin and the electron gives rise to the magnetic hyperfine splitting (this is described in many quantum mechanics textbooks). The form of the interaction

term in the atomic Hamiltonian is  $H_{hyp} \propto J \cdot I$ , which results in an energy splitting

$$\Delta E = \frac{C}{2} [F(F+1) - I(I+1) - J(J+1)]$$

where  $F = I + J$  is the total angular momentum quantum number including nuclear spin, and  $C$  is the “hyperfine structure constant.” Figures 7 and 8 shows the lower S and P energy levels for  $^{85}\text{Rb}$  and  $^{87}\text{Rb}$ , including the hyperfine splitting.

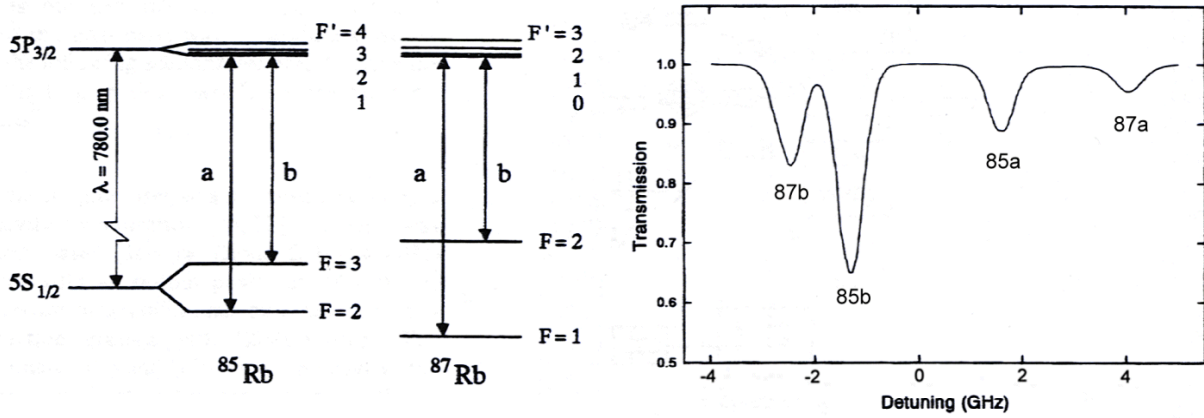


Figure 7. (Left) Level diagrams for the D2 lines of the two stable rubidium isotopes. (Right) Typical absorption spectrum for a rubidium vapor cell, with the different lines shown.

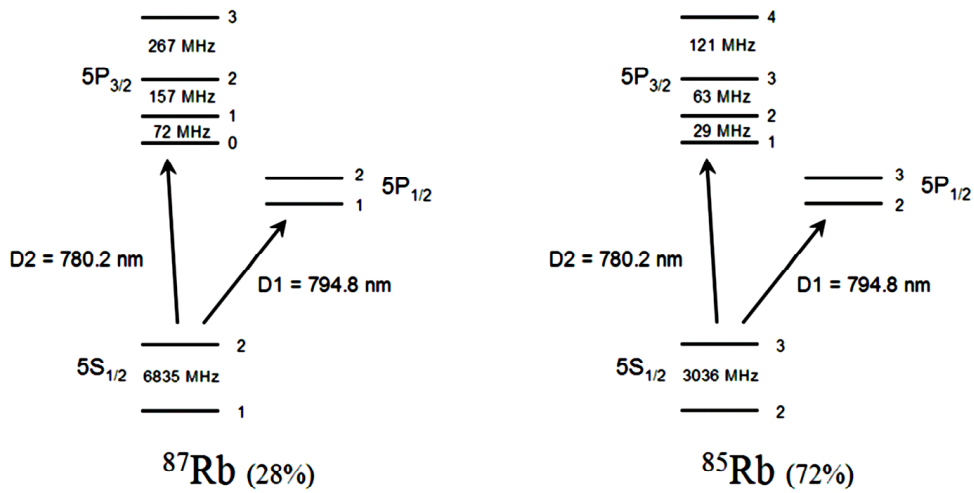


Figure 8. More rubidium level diagrams, showing the hyperfine splittings of the ground and excited states.

## II. LABORATORY EXERCISES.

The goal of this section is first to observe and record saturated absorption spectra for as many of the



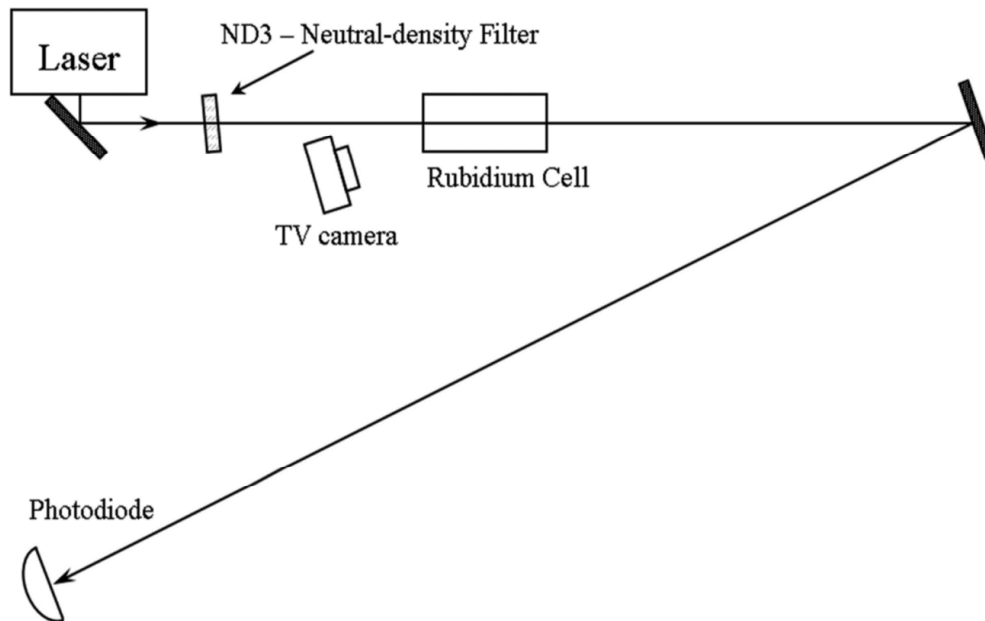


Figure 9. Recommended set-up to get the laser running on the rubidium resonance lines.

rubidium lines as you can, and then to see how well you can measure the  $P_{3/2}$  hyperfine splitting of  $^{87}\text{Rb}$  using a auxiliary interferometer as a length standard.

Remember that eye safety is important. First of all the laser operates at 780 nm, which is very close to being invisible. Thus you can shine a beam into your eye without noticing it. Also, the laser power is about 20 milliwatts, and all that power is concentrated in a narrow beam. Looking directly at the Sun puts about 1 milliwatt into your eye, and that much power is obviously painful. It is certainly possible to cause permanent eye damage using the Ph76 laser if you are not careful. Therefore – be careful. ALWAYS WEAR LASER GOGGLES WHEN THE LASER IS ON! As long as you keep the goggles on, your eyes will be protected.

### Week 1 – Getting the Laser On Resonance.

The first step is to get the laser turned on and tuned to hit the rubidium lines. We see in Figure 7 that the lines span about 8 GHz, which can be compared with the laser frequency of  $\nu = c/\lambda = 4 \times 10^{14}$  Hz. Thus to excite the atoms at all the laser frequency must be tuned to about a part in  $10^5$ . Start with the simple set-up shown in Figure 8. The ND filter can be removed when aligning the laser beam.

Once you have the beam going about where you want it, sweep the high-voltage going to the grating PZT with a triangle wave, so that the voltage varies from about 0 to 100 volts. Use the HV/100 to monitor the high voltage on the oscilloscope. Sweeping this voltage sweeps the grating position using a small piezo-electric actuator (made from lead zirconate titanate, hence PZT). While the high voltage is scanning you should then also change the laser injection current up and down by hand. The current makes large changes in the laser frequency, while the PZT makes small changes (see the laser primer for details).

The plan is that with all this sweeping the laser will sweep over the rubidium lines and you will see some fluorescence inside the vapor cell. This will appear as a bright line inside the cell; don't be confused

by scattering off the windows of the cell. If you cannot see the atoms flashing at all, ask your TA for help. The laser may need some realignment, or you may just not be doing something right.

Once you see fluorescence, compare the photodiode output to the rubidium spectrum shown in Figure 7. Usually you can only get the laser to scan over part of this spectrum without mode hopping (see the laser primer). Record your best spectrum using the digital oscilloscope and print it out. At this point the laser is tuned to the rubidium lines. Before proceeding with the rest of the experiment, move the ND filter from its location in Figure 9 to a new position right in front of the photodiode. If you look closely you'll see the absorption lines are still there, but much weaker. How come? There are two reasons. First, optical pumping is faster with more laser power, so the atoms are more quickly pumped to the dark state. That makes the absorption less. Second, the atoms become saturated with the high power, just like you calculated above. That also reduces the absorption.

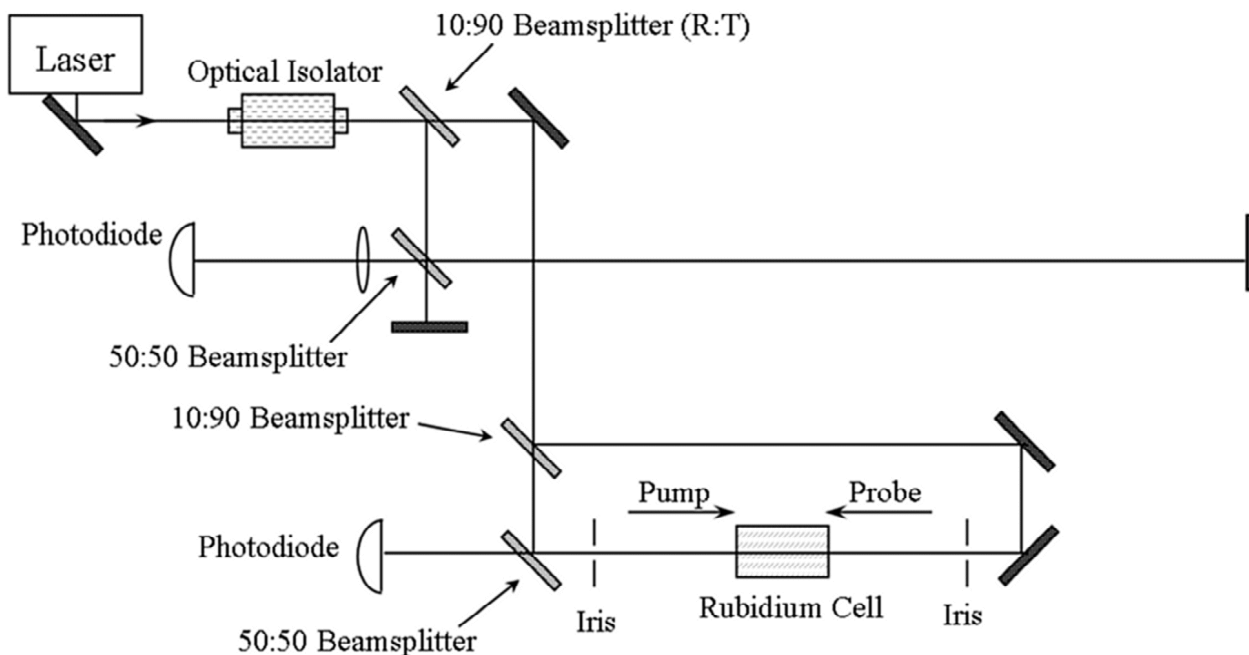


Figure 10. Recommended set-up to record rubidium saturated absorption spectra, and for measuring the hyperfine splittings.

### Week 1 – Getting a Saturated Absorption Spectrum.

The suggested set-up for observing saturated absorption spectra is shown in Figure 10. Since the laser is on resonance from the last section, leave it alone while you change the set-up. Ignore the interferometer part for now; that comes in after you've gotten some spectra. The optical isolator is a device that contains a two polarizers, a special crystal, and strong permanent magnets (see Appendix I). The first polarizer is aligned with the polarization of the input laser (vertical), and simply transmits the beam. The crystal in the magnetic field rotates the polarization of the beam by about 45 degrees, using the Faraday effect, and the beam exits through the second polarizer, which is set at 45 degrees. A beam coming back toward the

laser sees all this in reverse; the beam polarization gets rotated in the crystal, so that the polarization is 90 degrees with respect to the vertical polarizer, and the beam is not transmitted. These devices are also sometimes called optical diodes, since light only passes through them in one direction. We use an optical isolator here to keep stray light (generated downstream...note the pump beam goes backward after it passes through the cell) from getting back to the diode laser, where it can adversely affect the frequency stability.

Note the 10:90 beamsplitter puts most of the laser power into the probe beam. The irises are an alignment guide; if you have both the pump and probe beams going through small irises, then you can be assured that the beams overlap in the rubidium cell. If you block the pump beam you should get a spectrum that looks pretty much the same as you had in the previous section.

**Lab Exercise 1.** Observe and record the best spectra you can for whatever rubidium lines you can see, especially the two strongest lines (87b and 85b in Figure 7). Get some nice spectra and put hard copies into your notebook. Note (but don't bother recording) that the saturated absorption features go away if you block the pump beam, as expected.

### Week 2 – Measuring the Hyperfine Splitting.

Now finish the set-up in Figure 10 by adding the interferometer. (Turn off the laser frequency scanning while setting up the interferometer, so the fringes are stable.) Make the arm difference as long as you can. If you want you can add another mirror to the long arm to bounce it across the table. The longer the long arm, the better your measurement will be. Recombine the beams on the beamsplitter and send one of the output beams through a strong lens, so that the beam is expanded quite a bit. Align the overlap of the beams (in position as well as angle) until you see nice fringes on the expanded beam. Align the overlap so the fringes spacing is very large. Place the photodiode such that it only intercepts the light from one fringe of the interferometer.

If you now scan the laser frequency you should observe temporal fringes on the photodiode output. The fringe spacing can be computed from the arm length difference, which you should measure. When a beam travels a distance  $L$  it picks up a phase  $\varphi = 2\pi L/\lambda$ , so the electric field becomes

$$E = E_0 e^{i\omega t} e^{i\varphi}$$

When the beam is split in the interferometer, the two parts send down the two arms, and then recombined, the electric field is

$$\begin{aligned} E &= E_{arm1} + E_{arm2} \\ &= E_0 e^{i\omega t} \left[ e^{i4\pi L_1/\lambda} + e^{i4\pi L_2/\lambda} \right] \end{aligned}$$

where  $L_1$  and  $L_2$  are the two arm lengths. The additional factor of two comes from the fact that the beam goes down the arm and back again. Squaring this to get the intensity we have

$$\begin{aligned} I &\sim \left| e^{i4\pi L_1/\lambda} + e^{i4\pi L_2/\lambda} \right|^2 \\ &\sim 1 + \cos\left(\frac{4\pi\Delta L}{\lambda}\right) \end{aligned}$$

where  $\Delta L = L_1 - L_2$ . If the laser frequency is constant, then the fringe pattern goes through one cycle every time the arm length changes by  $\lambda/2$ .

**Problem 3.** If  $\Delta L$  is fixed, how much does the laser frequency have to change to send  $I$  through one brightness cycle? For your known  $\Delta L$ , what is the fringe period in MHz? From this you can convert your

measurement of  $\Delta L$  into a calibration of the laser frequency scan.

Use the two oscilloscope traces to plot the interferometer fringes and the saturated absorption spectra at the same time, as you scan the laser frequency. Watch that the interferometer fringes are uniform as a function of PZT voltage; if not the nonlinearities could compromise your calibration. Zoom in on the hyperfine features you want to measure. You will need to know which features belong to which lines, so identify the features by comparing your spectra with the level diagrams in Figures 7 and 8. Print out some good spectra, measure the spacings of the various features using a ruler, and you can turn this all into a direct measurement of the hyperfine splittings. Try to do this for both lines 87b and 85b in Figure 7. Note there are no tricks or complicated math in any of this. You just have to understand what's going on, and not lose any factors of two. No fair adjusting the answer by factors of two until it agrees with the known splittings.

**Lab Exercise 2.** Measure and record the largest  $P_{3/2}$  hyperfine splittings for  $^{85}\text{Rb}$  and  $^{87}\text{Rb}$ , in MHz. Estimate the accuracy of your measurement, knowing the various uncertainties you encountered along the way.

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## Appendix I – The Optical Isolator

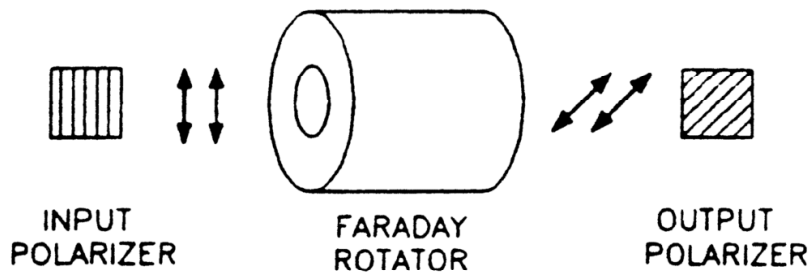


Figure 11. Schematic picture of an optical isolator. Not shown is the large longitudinal magnetic field in the Faraday rotator produced by strong permanent magnets inside the device.

The optical isolator is a somewhat subtle device, which uses the Faraday effect. The Faraday effect is a rotation of the plane of polarization of a light beam in the presence of a strong magnetic field along the propagation axis. You can get a feel for this effect by considering a simple classical picture. An incoming light beam imposes an oscillating electric field on the electrons in the solid, which causes the electrons to oscillate. Normally the oscillating electrons re-radiate the light in the same direction as the original beam, which doesn't change the polarization of the light (it does change the phase, however, which is the cause of the material index of refraction). With the application of a strong longitudinal magnetic field, you can see that the Lorentz force  $e \vec{v} \times \vec{B}$  will shift the motion of the electrons, and rotate their plane of oscillation. As the electrons re-radiate this tends to rotate the polarization of the light beam. Obviously a hand-wavy argument, but it gives you the right idea.

The optical isolator uses the Faraday effect to rotate the polarization angle of the input beam by 45 degrees, and the output beam exits through a 45-degree polarizer (see Figure 12). Note that the diode laser's beam is polarized, in our case along the vertical axis. If one reflects the beam back into the optical isolator, the polarization experiences another 45-degree rotation, in the same direction as the first, and the beam is then extinguished by the input polarizer. You can see that the rotations have the correct sense using the classical picture. Thus the overall effect is that of an "optical diode" – light can go through in one direction, but not in the reverse direction.

The Faraday effect is typically very weak, so the optical isolator uses a special crystal, which exhibits an anomalously large Faraday effect, and a very strong longitudinal magnetic field produced by state-of-the-art permanent magnets. Optical isolators have gotten much smaller over the last couple of decades as magnet technology has improved. The magnetic field is strong only near the axis of the device, which therefore has a small clear aperture. Also, too much light intensity will burn a spot in the Faraday crystal, so one must be careful not to focus the diode laser to a tight spot inside the optical isolator.